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Structure of Accretion Shocks and Radiative Cooling Effects in High Energy Density Plasma Experiments

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the Diploma of Imperial College London

Declaration of Originality

I hereby certify that all work contained within this thesis is entirely original and my own work, unless otherwise stated.

Stefano Merlini

October 23, 2023

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Abstract

Accretion shocks are ubiquitous phenomena in many astrophysical systems which can be significantly affected by radiative cooling effects, leading to instability and turbulence. This thesis presents a study of accretion shock experiments at the MAGPIE pulsed power facility. Two methods are used to produce reverse shocks in laboratory. The first investigates the structure of reverse shocks resulting from the collision of supersonic, magnetised plasma flows generated by an inverse wire array interacting with a planar conducting obstacle, whereas the second approach consists of the ablation by X-ray radiation from a Z-pinch wire array of solid targets. In the planar obstacle experiments, variations in the reverse shock structure are observed depending on the wire material used, despite similar upstream flow velocities and mass densities. Specifically, when aluminium wire arrays are employed, a well-defined, sharp shock is formed that aligns with magneto-hydrodynamic theory. However, in the case of tungsten wires, a distinct stand-off shock is not observed, instead, a broad region ahead of the obstacle is characterised by density fluctuations spanning a wide range of spatial scales. These two contrasty interactions are diagnosed with laser interferometry, Thomson scattering, shadowgraphy, and a newly developed imaging refractometer, enabling to characterise the small-scale density perturbations by detecting deflections of the probing laser. These measurements suggest that the differences in shock structure are most likely due to radiative cooling effects which give rise to density perturbations elongated along magnetic field lines. In the second experimental campaign, reverse shocks formed from the collision of counter-streaming supersonic plasma flows are investigated. The properties of the shocked layer were determined using various laser diagnostics and optical self-emission images, revealing overall consistency with a 1-D accretion shock model for a value of $\gamma \leq 1.2$. In addition, the capability of the X-ray-driven platform to easily change the B-field orientation allowed us to explore the case for colliding flows in a transverse magnetic field, revealing promising results for investigating new astrophysical systems in the laboratory.

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Chapter 1

Introduction

Plasma, a mixture of electrically charged particles, constitutes the majority of matter in the visible universe. This state of matter exists in a wide range of densities and temperatures, from the tenuous and cold interstellar medium to the dense and hot centre of giant supernovas formed from the collapsing of stars. Plasmas are ionised gases which can be profoundly affected by the presence of magnetic fields, determining their evolution and dynamics. There are countless plasma phenomena in nature mediated by the presence of magnetic fields, and one of the most important processes among them is the magnetospheric accretion in young stellar objects (YSO).

The current model of star formation consists of four main phases. First, massive clouds of dust and gas (mainly hydrogen), also known as nebulae, start to be attracted together by their gravity. Hydrogen gas starts to spiral around the nebula core, heating up and eventually transitioning into a protostar. At this early phase, the protostar still gathers mass from its parent molecular cloud, and fusion reactions are not yet self-sustained. As the protostar grows and incorporates almost all the surrounding interstellar material, the protostar then contracts inducing the rise of the internal temperature and fusion reactions begin. This is the start of the pre-main sequence stage, where stars at this phase are known as young stellar objects and they can be classified based on their mass as T Tauri stars [1, 2] for low mass YSO ($M < 2M_{\odot}$), or Herbig Ae/Be stars [3] for high mass pre-main sequence objects ($2M_{\odot} < M < 8M_{\odot}$).¹

The accretion dynamics in the formation of young stars is still a matter of debate because of the

¹ $M_{\odot}$ is the solar mass, approximately 2×10^{30} Kg

lack of direct observations and modelling; however, the most accredited mass accretion mechanism for both T Tauri and Herbig Ae/Be stars is the magnetospheric accretion. Originally proposed by [4] from the previous work by [5, 6], magnetospheric accretions consist of matter from the star surrounding (temperature, $T_1 \simeq 2000$ K and electron density, $n_e = 10^{11} \div 10^{13} \text{ cm}^{-3}$) being accreted onto the star surface through accretion columns mediated by the stellar magnetic field. These accretion channels connect the star surface to the edges of the external disk with the magnetic field lines guiding the plasma material at supersonic free-fall velocity ($u_1 \sim 500$ km/s) towards the chromosphere. The impact of the plasma onto the stellar surface produces a reverse shock, which expands back into the accretion stream as more material is heated ($T_2 \simeq 10^6$ K) in the post-shock region with significantly reduced velocity ($u_2 \sim 100$ km/s). While there is clear evidence of magnetospheric accretion columns in T Tauri stars, mainly thanks to strong magnetic fields [7], observations in Herbig Ae/Be objects are more ambiguous. The magnetic pressure in these stars is too weak to truncate the surrounding accretion disk, resulting in only partial gaps between the stars and their disks [8–10]. Another category of stellar objects where magnetospheric accretion occurs is magnetic cataclysmic variables (MCV). These astrophysical objects are semi-detached binary systems containing a highly magnetised white dwarf accreting material coming from its Roche-lobe-filling companion [11, 12]. Within this category, systems named AM Her stars or polars are the most relevant for this work, as the magnetic fields of the white dwarf are strong enough to prevent the formation of an accretion disc around the star. Similar to YSO, plasma material leaving the companion and tied to the field lines of the white dwarf is led toward the magnetic poles of the compact star where accretion columns form.

In both cases, characteristics of these systems, such as mass accretion rates, are typically inferred through the analysis of direct optical observations [13] by deducing the star-accretion luminosity from its optical emission lines. More recently, the mass accretion rate of classical T Tauri stars has been derived from their soft X-ray emission [14, 15]. In all the mentioned methods, it is assumed that the stream impacts the stellar surface with free-fall velocity from a distance equal to a few times the radius of the star, and the plasma conditions of the post-shock region are stationary over time. These simplistic assumptions and discrepancies between the observed X-ray luminosity and the predictions based on the UV-visible bands from the same object indicate that there are still unsolved challenges. Moreover, the lack of reliable measurements limits our current capabilities to numerically study the full

dynamics of the magnetospheric accretion column. Current numerical magnetohydrodynamic models are restricted to two spatial dimensions with limited access to radiative processes such as cooling instabilities, which are known to play an important role in the dynamics of accretion shocks [16].

An alternative approach that can complement astrophysical observation and theoretical study is a laboratory of plasma physics where detailed measurements of plasma properties and evolution can be made in scaled systems of astrophysical relevance. To produce astrophysical-relevant plasma experiments in the laboratory, high pressure is typically required, and the field is often referred to as High Energy Density Laboratory Astrophysics (HEDLA). This is usually achieved using high-energy lasers or pulsed power generators, which can generate plasma conditions that are scalable to astrophysical objects. The set of equations, the MHD equations, that govern the dynamic of the experiments are the same as those governing the astrophysical systems, and as long as scaling criteria are satisfied, the two systems will behave similarly, allowing to investigate with great detail the underlying physics. Details are not given here, but can be found in [17–19]. In the last decade, a few experiments have been designed to replicate accretion shocks and all of them were laser-driven experiments [20–22]. This thesis aims to expand further the available designs using the pulsed-power generator MAGPIE as the driver to replicate accretion flows in a highly controllable environment, improving our current understanding of these accretion mechanisms responsible for the formation of new stars.

1.1 Theoretical Model of Accretion Shocks

The impact of the accretion stream onto the stellar chromosphere generates reverse shocks which propagate upstream through the accretion column creating a post-shocked region of hot plasma. In its simplest form, a shock is formed when supersonic plasmas collide against an obstacle, forming an abrupt change in plasma parameters and an increase in entropy. Material is not only compressed at the shock, it is also heated. The stronger the shock, the denser and hotter will be the shocked material. In classic unmagnetised shocks, this heating process is mediated by ion viscous effects through ion-ion Coulomb collisions. In this case, the transition where dissipation mechanisms occur is typically very

small and not resolvable in both experiments and numerical simulations, therefore this shock transition is often approximated to a hard boundary with negligible extent between the unshocked upstream flow and the post-shocked plasma. Although entropy is net positive in shock systems, where kinetic energy is converted into thermal motion, conservation laws can still be applied across the shock as long as both upstream and downstream conditions are considered far enough from the shock front. Under this condition, the properties of the shock system are described by the shock jump conditions. A simple schematic representation of a one-dimensional shock is shown in Fig. 1.1. For the sake of brevity, we restrict the description only to a few useful relations that will be used in this thesis, leaving their derivations to the available textbooks [23–25].

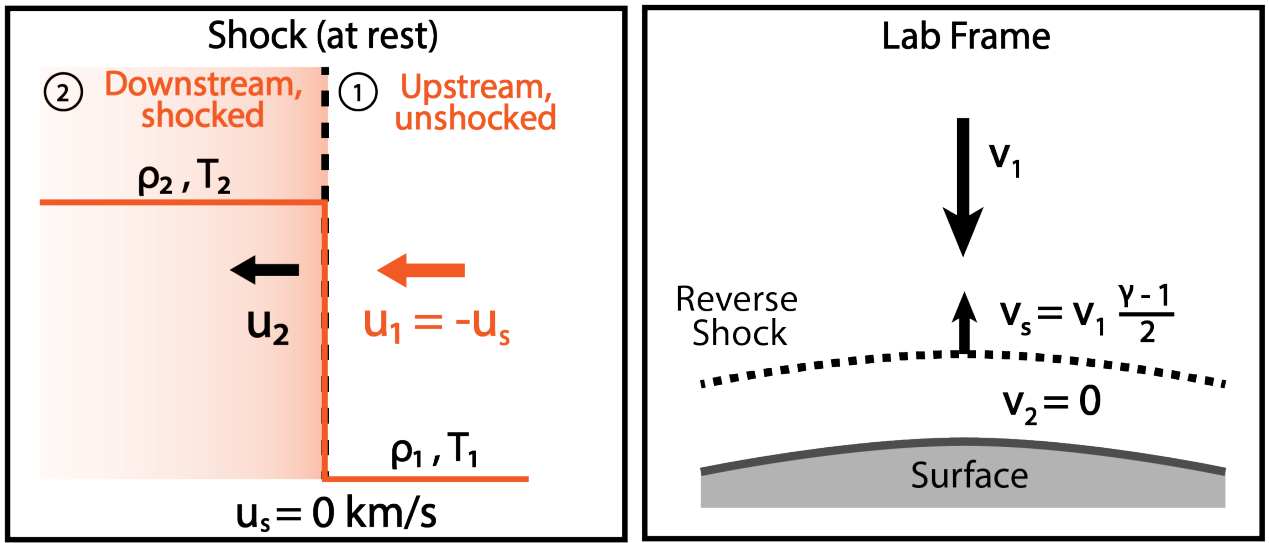


Figure 1.1: (Left) Schematic illustration of a one-dimensional shock in the shock frame of reference. The shock is at rest ($u_s = 0 \text{ km/s}$) with the post-shock flow leaving the shock front with velocity $u_2 < u_1$ and the unshocked upstream material flows into the stationary shock from the right with velocity $u_1 = -u_s$. (Right) In the lab frame, the post-shock shock region is at rest, $v_2 = 0 \text{ km/s}$ with the shock moving away from the rigid surface with shock speed v_s .

In the shock frame of reference, the shock front is at rest $\mathbf{u}_s = 0 \text{ km/s}$ with the incident unshocked flow (subscript 1) moving at velocity $\mathbf{u}_1$ while the flow in the post-shock region (subscript 2) leaves the shock front at speed $\mathbf{u}_2 < \mathbf{u}_1$. The shock jump relations are derived in this frame of reference by imposing the three fluid equations (the Euler equations) describing the conservation of mass, momentum and energy:

$$\rho_1 v_1 = \rho_2 v_2 \quad (1.1)$$

$$\rho_1 u_1^2 + p_1 = \rho_2 u_2^2 + p_2, \quad (1.2)$$

$$\rho_1 u_1 \left(\varepsilon_1 + \frac{u_1^2}{2} \right) + p_1 u_1 = \rho_2 u_2 \left(\varepsilon_2 + \frac{u_2^2}{2} \right) + p_2 u_2 \quad (1.3)$$

Where ε is the internal energy per particle defined by the equation of state (EOS) of the plasma. These jump conditions can be used to obtain useful equations for density compression, velocity jump and temperature increase across the shock front. Assuming an ideal gas with polytropic index γ , the EOS is given by $p\varepsilon = p/(\gamma - 1)$ which leads to the pressure in the post-shock region:

$$p_2 = \frac{2\rho_1 u_s^2}{\gamma + 1} \left[1 - \frac{(\gamma - 1)p_1}{2\rho_1 u_s^2} \right]$$

For very fast shock, most of the kinetic energy carried by the upstream ions is efficiently converted into heat as the pressure behind the shock increases as $p_2 \propto u_s^2$. This class of shocks are known as strong shocks characterised by an upstream sonic Mach number $M_u \gg 1$. The upstream Mach number is defined as the ratio between the upstream velocity and the sound speed of the unshocked flow, $M_u = u_1/c_{s,1}$. This condition is typically the one considered in accretion shocks, where plasma material falls under the gravitation attraction of the star at high supersonic speed. In this shock regime, the compression of mass through the shock is solely dependent upon the polytropic index with the velocity ratio exactly opposite and equal to the density jump:

$$\frac{\rho_2}{\rho_1} = \frac{u_1}{u_2} = \frac{\gamma + 1}{\gamma - 1}$$

It is important to note that the value of γ depends on the number of degrees of freedom of each particle in the plasma. In an ideal monoatomic gas, energy can be distributed within the system only by three translation degrees of freedom per atom, $\gamma = 5/3$, while in the case of the radiative dominant regime, the value of γ is usually much smaller, sometimes approaching unity. This is the case for radiative shocks, where strong compression is often observed, leading to post-shock regions that are much colder and dense than standard shocks. In the context of accretion shocks, this effect can have a strong impact on the thickness of the shocked slab and the post-shock temperature, ultimately affecting the amount of material contributing to accreting the mass of the star. An expression of the post-shock temperature, T_2 , can be obtained by rewriting the eq. 1.1 in the strong shock limit and substituting $p_2 = (Z_2 + 1)k_B \rho_2 T_2 / Am_p$ where Z_2 and Am_p are the ionisation state of the plasma in the post-shock

region and the atomic mass, respectively. If any ionisation processes are neglected across the shock, the final temperature in the shocked region is:

$$k_B T_2 = A m_p u_s^2 \frac{2(\gamma - 1)}{(\gamma + 1)^2}$$

This description of the one-dimensional shock system presented above is typically the reference model for reverse shocks where a hard surface offers the supporting pressure to the stagnant shocked plasma and similarly, it is the one found in the literature of astrophysical accretion shocks. It is worth noting that the 1-D reverse shock model is a good approximation of magnetospheric accretions only if the stellar magnetic field is uniform and locally strong, $\beta = P_{th,2}/P_B \ll 1$. In the case of a post-shock region characterised by high- β , the 1-D approximation is no longer justifiable as shocked plasma can escape laterally at the surface star or bounce back into the accretion stream, invalidating any possible predictions on the shock position, hence mass accretion rate. In those cases, 2D or 3D radiative MHD shock models are necessary [26, 27].

Until now, the expressions presented above were derived in the shock frame of reference considering the upstream velocity, $\mathbf{u}_1$. In laboratory experiments, however, this value is often not available, especially when the post-shock plasma is not moving supported by the pressure of an unmovable surface. Moreover, all the measurements are performed in the laboratory frame, therefore it is useful to provide the relevant relations between the two frames of references, as shown in Fig. 1.1. Under the hypothesis of strong shocks, the velocities in the lab frame are obtained by subtracting $\mathbf{u}_2$ so that $\mathbf{v}_1 = \mathbf{u}_1 - \mathbf{u}_2$ and the shock propagates at the velocity:

$$v_s = -v_1 \frac{\gamma - 1}{2}$$

In this simple 1D accretion shock model, radiative effects are reflected only in the choice of the polytropic index γ , ignoring any potential instabilities on the shock front associated with rapid cooling of the shocked region. In this work, we aim to go beyond this simplistic description of a reverse shock including radiative phenomena that can potentially destabilise the shock until the development of turbulence. In the following section, radiative cooling instabilities are introduced together with a practical framework for those relevant to the experiments.

1.2 Radiative cooling instabilities

Radiative cooling instabilities are often referred to as the broad range of cooling effects that can occur when the plasma rapidly loses heat energy through radiation. These are often very common in astronomy, specifically in extreme phenomena when plasma material undergoes rapid heating and cooling in a short timescale. This is particularly true in the magnetospheric accretion process of young stellar objects, where funnelled streams of plasma originating from the surrounding disk of the protostar fall at supersonic velocity against the stellar surface. The impact on the chromosphere generates strong reverse shocks that heat the inflow of materials at a few million Kelvin. These shocks are naturally radiative and their structure is strongly correlated with the characteristic cooling time at which they operate ($\sim 10 - 10^3$ s). Under these thermal pressures in the post-shock region and for optically thin plasma, the reverse shocks are launched along the accretion streams, before the hot slab radiates its energy away and collapses. This process is expected to repeat quasi-periodically, which results in pulsating X-ray emission, as has been demonstrated by 1D numerical studies [28–30]. Several direct observations of accreting stellar systems, however, revealed the absence of these periodic oscillations in the X-ray spectral emission of these objects [31, 32], leading to open questions on the stability of the reverse shock and the contributions of other radiative effects. Cooling instabilities, together with extra spatial dimensions, can lead to inhomogeneous reverse shocks with no periodicity [27, 33] or it is arguably possible that these radiative effects stabilise the reverse shock inside the chromosphere of the star. In both cases, the role of radiative cooling effects plays a crucial role in the systems and can significantly compromise the mass accretion of young stars. The present work attempts to explore these radiative effects in the laboratory, replicating reverse shocks where radiation losses are dominant. In general, radiative cooling instabilities can be classified into three types: shock front oscillation, thin-shell instability and thermal instability. A schematic illustration is shown in Fig. 1.2.

The first one is better known from the work of [16] and is the oscillatory behaviour theorised to occur in accretion shocks, just mentioned above. This is also one of the few instabilities that occurs in 1-D systems. The second type, known as Thin-Shell instability or shock-bounded slab instability, occurs along shocks when a dense narrow region with an internal pressure is formed behind the shock front, supported by a constant pressure at its rear edge. This narrow and dense 2-D post-shock region

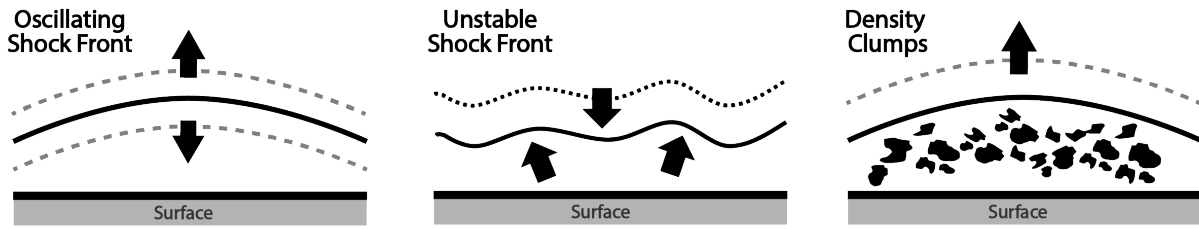


Figure 1.2: Schematic illustration of three different types of cooling instability. The first one on the left is one of the few instabilities that can occur in a 1-D system and is mostly known from [16]. The remaining two are usually referred to as thin-shell instability [34] and Thermal instability [35], respectively, occurring in 2-D or 3-D space.

can form in both decelerating or accelerating shocks in a strong radiative cooling regime. The case of accelerating shocks can be thought of as a compressible, thin-shell variant of Rayleigh -Taylor instability [34, 36, 37]. This is typically common in counter-streaming colliding flows rather than accretion shocks; however, there are no obvious reasons why reverse shocks should not be as well susceptible to thin-shell instability as moving radiative shocks under a continuous inflow of materials. The third and last type of instability is known as thermal instability [35]. This is an uncontrolled runaway process, leading to characteristic density clumps and controlled at small scales by thermal conduction. In contrast to the first instability, both the thin-shell instability (TSI) and thermal instability can only occur in 2-D or 3-D space. In this thesis, we will only focus on these two phenomena as relevant to the experiments presented later while Chevalier & Imamura instability will not be discussed.

1.2.1 A simple Thin-Shell Instability Model

Thin-shell instability is a common phenomenon that occurs in many astrophysical shocks [38, 39] and laboratory experiments [40, 41] involving thin layers of matter that lose energy through radiation. These layers tend to grow in mass as they move, but they can become unstable in the presence of pressure imbalance between the inflowing upstream plasma and the internal pressure acting always normal to the layer surface. This instability was first proposed by [42] and confirmed experimentally by [43]. In later studies, this instability is sometimes called the Vishniac-Ryu instability [34] for shocks that slow down, or the “Thin-Shell instability” and the “nonlinear Thin-Shell instability” for colliding flows [36, 44, 45]. In general, a radiative reverse shock can be considered as a two-dimensional thin-shell layer that accumulates mass from the unshocked plasma as it propagates into the accretion

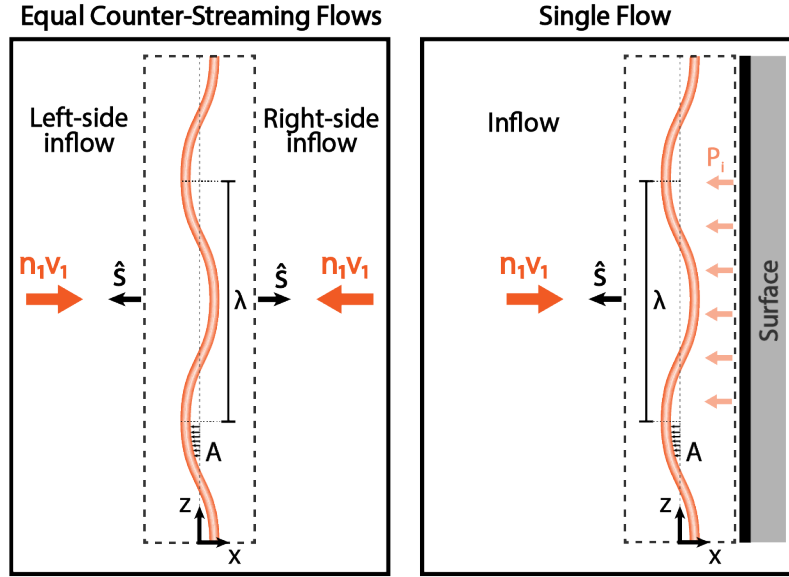


Figure 1.3: Schematic illustration of the Thin-Shell instability model. The model is originally derived considering equal counter-streaming flows, however, the same arguments can still be made for an accretion shock driven by a constant inflow, $n_1 v_1$ and supported by a constant pressure, P_i .

stream. This thin layer is bounded from one side by the shock front and from the other side by the supporting pressure of the shocked plasma. In the strong shock limit presented in the previous section, the thickness of the layer is expected to grow linearly in time, $\xi(t) \sim c_{s,2}^2 t / u_1$, as a function of the sound speed $c_{s,2}$ defined in the post-shocked region and the velocity u_1 of the incoming flow [36]. In case of small perturbations along the layer, the denser parts of the layer slow down more than the less dense parts, reversing the perturbation and causing a non-homogeneous growth either faster or slower of the layer, which can lead to the disruption of the shock.

A simple model to describe the growth of sinusoidal instabilities on thin-shell structures was introduced by [45] assuming a thin-shell layer formed by two equal and opposite colliding flows. Although the model is not derived directly from the typical geometry of accretion shocks, the same arguments can still be applied considering only the ram pressure of a single plasma inflow. A summary of the model is presented below, where an expression of the growth amplitude of a single mode, λ , is derived by simply imposing conservation laws across the layer.

The system under consideration is represented in Fig. 1.3. The collision of two counter-streaming flows with equal ram pressure gives rise to the formation of a thin-shell layer in volume V with S (indefinitely extended along the z -axis) its surface contour and $\hat{S}$ the vector normal opposite to inflow

directions. In this model, the shock is at rest with $v_s \sim 0$ km/s. In the experiments, this is a valid approximation as long as $v_s \ll v_1$. Any variations of mass and momentum within the volume V must be equal to respective fluxes of mass and momentum crossing the surface contour S of the volume, according to the fulfilment of the conservation laws:

$$\Delta N_i = N_i(t) - N_i(t_0) = 2n_1 v_1 \lambda (t - t_0)$$

$$\Delta \mathbf{p} = \mathbf{p}(t) - \mathbf{p}(t_0) = 0$$

where ΔN_i and $\Delta \mathbf{p}$ are the variations of the total number of ions reaching the layer and the variation of the linear momentum per unit length in volume V , respectively. Because we are in the context of shocks, only the ions contribute significantly to the overall dynamics of the thin-shell layer. Moreover, under the hypothesis of perfectly collimated flows, any fluxes leaving the sides of the volume are considered negligible. As a result, any changes in ion flux that reach the layer at a given time, $N_i(t)$, must be associated with an equal change in layer mass:

$$\rho_i(t) L \xi(t) \approx N_i(t) = N_i(t_0) + n_0 v_0 \lambda (t - t_0)$$

with $\rho_i(t)$ the layer mass density, $\xi(t)$ the layer thickness and L is the total length of the layer. If we assume an initial sinusoidal perturbation of wavelength λ along the x -axis, $f(x, t) = A(t) \sin\left(\frac{2\pi}{\lambda}x\right)$ starting at $t_0 = t_{0,\lambda_0}$, the total length of the thin shell will be given by the arc length integral:

$$L = \int_0^\lambda \sqrt{1 + \left(\frac{\partial f(x, t)}{\partial x}\right)^2} dx \rightarrow L(A(t)) = \lambda + \frac{\pi^2 A(t)^2}{\lambda} + \dots$$

Solving the integral and expanding in series the final expression about $A(t_{0,\lambda_0}) = 0$, we obtain:

$$\rho_i(t) \left[\lambda + \frac{\pi^2 A(t)^2}{\lambda} \right] \xi(t) = N_i(t_{0,\lambda_0}) + 2n_0 v_0 \lambda (t - t_{0,\lambda_0})$$

At the time $t = t_{0,\lambda_0}$, we are at the beginning of the perturbation, $N_i(t_{0,\lambda_0}) = \rho_{(i,0)} \xi_0 \lambda$ where $\rho_{(i,0)}$ and

ξ_0 are the initial ion mass density and thickness of the layer.

$$A(t) = \frac{\lambda}{\pi} \sqrt{\frac{2n_0 v_0 (t - t_{0,\lambda_0}) + [\rho_{(i,0)} \xi_0 - \rho_i(t) \xi(t)]}{\rho_i(t) \xi(t)}}$$

This expression of the perturbation amplitude is valid under the hypothesis of small sinusoidal perturbations with $A(t) \ll L$ (linear regime) and thin layers $\xi(t) \ll \lambda \leq L$ at the formation stage. It is important to note that in the case of a single driving ram pressure, the symmetry of the system is broken and, in general, the conservation of momentum, eq. 1.2.1, is non-zero. This can still be achieved by including other terms in the system, leading to a more complicated model described in [36]. Nonetheless, for the experiments presented in this thesis, the amplitude growth can still be used in the context of accretion flows by simply dropping the factor of 2 in the inflow flux of eq. (1.2.1).

1.2.2 Thermal Instability

Thermal instability is known to play a significant role in the evolution and formation of many astrophysical objects. Those objects cover a wide range of situations from solar prominences [46, 47] to interstellar media where self-gravitation alone cannot explain their existence as their calculated gravitational energies are far smaller than those due to internal pressure. Early studies from [48] and [49], proposed that those objects could be formed by condensation phenomena happening in the diffuse media due to unstable thermal equilibrium. Only years later, the correct instability criteria were derived [35], which still represents the widely accepted description of the problem [50]. When the rate of cooling due to radiation reaches a threshold, plasmas collapse, becoming unstable following what is called the radiative collapse phenomenon. Localised radiative cooling of the plasma can cause the temperature to drop, which leads to compression of the cooler volume as a result of a decrease in thermal pressure. This compression raises the density and therefore the cooling rate. This process is self-reinforcing, limited at small length scales by thermal conduction and at large length scales by the hydrodynamic motion of the plasma. In this section, the dispersion relation of the thermal instability for the unmagnetised case is provided along with a discussion of the instability criteria. The reader should refer to [35] for more details.

The analysis of the thermal instability starts assuming a non-expanding homogeneous medium in thermal equilibrium with no magnetic field. The basic governing equations are represented by the basic equations (Euler description) of mass, momentum and energy, considering thermal conductivity, k and heat losses, Q :

$$\frac{d\rho}{dt} + \rho \nabla \cdot \mathbf{v} = 0 \quad (1.4)$$

$$\rho \frac{d\mathbf{v}}{dt} + \nabla p = 0 \quad (1.5)$$

$$\frac{1}{\gamma - 1} \frac{dp}{dt} - \frac{\gamma}{\gamma - 1} \frac{p}{\rho} \frac{d\rho}{dt} - \rho Q - \nabla \cdot (k \nabla T) \quad (1.6)$$

where d/dt represents the convective derivative, $\partial/\partial t + \mathbf{v} \cdot \nabla$ and the energy equation is calculated by assuming an ideal gas with an arbitrary ratio of specific heat, γ . In the hypothesis of optically thin plasma, the thermal energy losses per unit of mass per second are simply a function of the local values of density, ρ and temperature, T , neglecting any radiative transfer effects. Under this condition, thermal equilibrium is achieved when $Q(\rho_0, T_0) = 0$.

If we now introduce perturbations in density and temperature such that a specific hydrodynamic variable remains constant (density, pressure and entropy), these will result in a change in thermal energy losses, δQ hence an associated change in system entropy, δS , according to the second law of thermodynamic, $\delta Q dt = -T d(\delta S)$. This means that the system becomes unstable for a constant hydrodynamic variable, A , only if the rate of change in heat-loss with respect to the system entropy S is positive:

$$\left(\frac{\partial Q}{\partial S} \right)_A > 0 \implies \left(\frac{\partial Q}{\partial T} \right)_A < 0 \quad (1.7)$$

Assuming exponential perturbations under the linear regime, the following cubic dispersion relation for the thermal instability can be obtained as described in [35]:

$$n^3 + n^2 c_s \left(k_T + \frac{k^2}{k_K} \right) + n c_s^2 k^2 + \frac{c_s^3 k^2}{\gamma} \left(k_T - k_\rho + \frac{k^2}{k_K} \right) = 0 \quad (1.8)$$

where n is the growth rate of the unstable modes with isochoric and isothermal wavenumbers k_ρ and k_T , both dependent upon the sign of Q , k_K represents the reciprocal of the mean free path of the conducting particles and c_s is the sound speed defined as $c_s = \sqrt{\gamma p_0 / \rho_0}$. According to eq. (1.8) there are three

possible solutions, depending on the signs of the temperature and density derivatives of the heat-loss function, Q :

- Condensation mode: this is a mode in which density and temperature variations occur with opposite phases so that pressure is constant. It satisfies the isobaric instability criteria $(\partial Q / \partial T)_p < 0$, which means that the cooling rate increases with decreasing temperature and increasing density. The growth rate decreases with increasing k and vanishes at a critical wavenumber k_c , which depends on the thermal conductivity of the gas.
- Wave mode: this case occurs when density and temperature variations have the same phase so that pressure variations are significant at constant entropy. This mode corresponds to sound waves, which are unstable $(\partial Q / \partial T)_s < 0$, meaning that the cooling rate increases with decreasing density. Similar to the condensation mode, the growth rate also decreases with increasing k and vanishes at the critical wave number, k_c , due to thermal conduction.
- Neutral mode: this mode represents the stable thermal equilibrium, where there is no growth or decay of perturbations for any wavenumbers, k . This mode occurs when neither of the instability criteria for the condensation or wave modes are satisfied.

A more practical stability criterion for thermal instability is provided by [51]:

$$\frac{d(\log \Lambda)}{d(\log T)} < C$$

where C is a constant that defines various thresholds of instability based on the three possible regimes: isobaric ($C = 2$), isochoric ($C = 1$) and adiabatic ($C = -0.5$). In astrophysical applications, the condition of instability is strongly dependent on the exact definition of the cooling function which can assume various forms based on the context. In its classic definition [51] the cooling function Λ is defined as $[\text{erg cm}^3 \text{ g}^{-2} \text{ s}^{-1}]$ so that $Q = \rho \Lambda$ is in units of $[\text{erg g}^{-1} \text{ s}^{-1}]$ with the mass density, ρ $[\text{g cm}^{-3}]$. In contrast to more recent literature where the cooling function appears as Λ_{ap} , which is expressed instead in units $[\text{erg cm}^3 \text{ s}^{-1}]$. For the instability analysis presented in this thesis, the cooling functions calculated from the FLYCHK code [52] will be used. FLYCHK returns a Total Power Loss in units of

($\text{erg s}^{-1} \text{ atom}^{-2}$), which is equivalent to the classic definition, Λ , after counting for the mass density of the element:

$$\Lambda_{FLYCHK} = \frac{\text{Total Power Loss}}{\rho A m_p} = \Lambda$$

An important parameter in the thermal instability analysis is represented by the cooling time formally defined as the ratio of thermal energy density over the radiative power density, $\tau_{cool} = U_{th}/P_{rad}$. This timescale represents the characteristic growth time of density perturbations due to thermal instability at near pressure equilibrium with its surroundings. For the cooling function Λ in [$\text{erg cm}^3 \text{ g}^{-2} \text{ s}^{-1}$], the expression for $P_{rad} = \rho^2 \Lambda$, while for the cooling function Λ_{ap} , this corresponds to $P_{rad} = n^2 \Lambda_{ap}$. In both definitions, the density dependence cancels in the stability criteria, and only the temperature dependence is important.

1.3 Experimental and Numerical Studies of Accretion Shocks

The development of high-powered lasers has led to a growing interest in high energy density laboratory of astrophysics, in which the dimensionless parameters of the plasma are scalable to study astrophysical objects [18, 19]. This includes accretion shock systems, where laser-driven experiments offer a potential path to investigate details of these accretion columns which are not resolvable directly from astrophysical observations. Two methods have been used to replicate accretion shocks in laser-driven experiments. The first approach generates plasma from laser ablation of solid targets and uses a physical tube to collimate the accreting flow, while the second and more recent designs involve the use of strong magnetic coils to generate the required magnetic pressure to collimate the plasma flow. The latter represents a more realistic representation of the astrophysical cases as in stellar accretion shocks, plasma material is brought towards the star surface by the magnetic field.

In early studies, the POLAR project was an attempt to replicate and investigate accretion shock structure under very strong radiative cooling regimes, using high-Z elements and lasers facility [53–55]. One of the limiting factors in their experiment was the duration of the plasmas (tens of nanoseconds) and the limited free access to the reverse shock with diagnostics. In contrast to the naturally occurring accretion columns in dwarf stars where plasma materials are collimated by the strong magnetic field

of the star towards the surface, in the experiment the collimation was instead achieved through a transparent tube. A schematic illustration is shown in Fig. 1.4(a), where plasma flows are formed from a shocked foam and launched into a vacuum tube towards the planar obstacle. The experiment revealed consistency between the growth of the radiative accretion shocks and the numerical predictions.

More recent experiments have been performed using laser platforms and external ambient magnetic fields [20–22] to study accretion columns in a more realistic setup allowing to investigate also the interaction of the lateral transverse post-shocked plasma with the ambient field. Similar to the earlier experiments, a laser is focused on a solid target launching a supersonic plasma plume which then collides against a flat obstacle, forming the reverse shock structure, as shown in Fig. 1.4(b). The

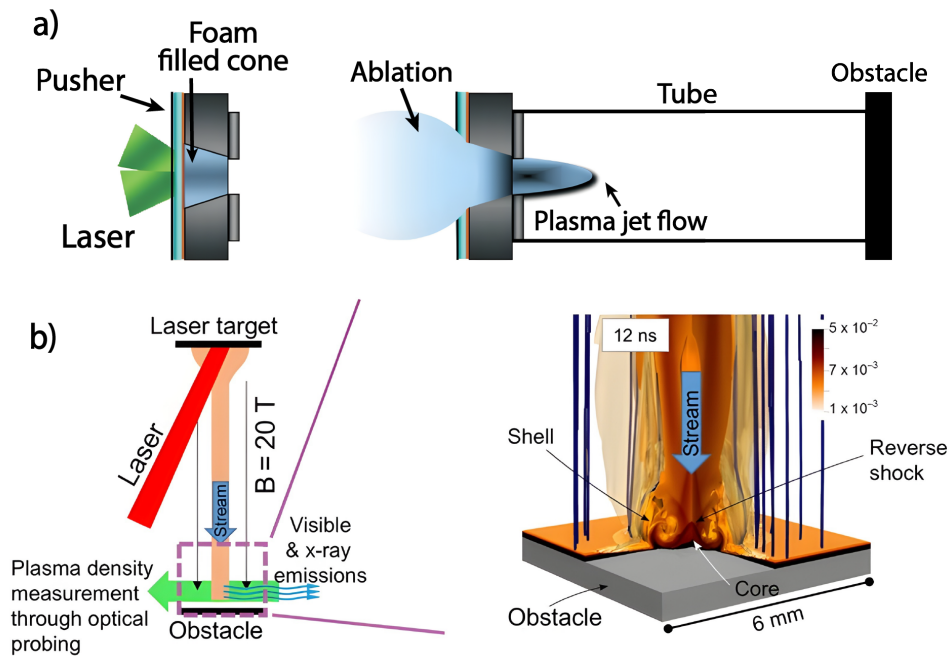


Figure 1.4: Schematic of laser-driven accretion shock experiments. (a) Plasma jets are launched by the laser impacts on a pusher and collimated by a physical transparent tube. Modified with permission from [53]. (b) A laser-driven experiment in which plasma flows are collimated by the presence of an external magnetic field parallel to the flow direction. Modified with permission from [21].

key difference here is the absence of physical boundaries, while still achieving collimation using an external magnetic field. These magnetically laser-driven experiments have been able to investigate the overall structure of the reverse shock, including the importance of the magnetic field strength on the X-rays emitted by the shocked core. X-ray emissivity profiles originating from young stellar objects are the only measurable parameters in such objects, playing the primary contribution in the estimates of the mass accretion rate. The experiments conducted by [21] revealed the complex 3-dimensional

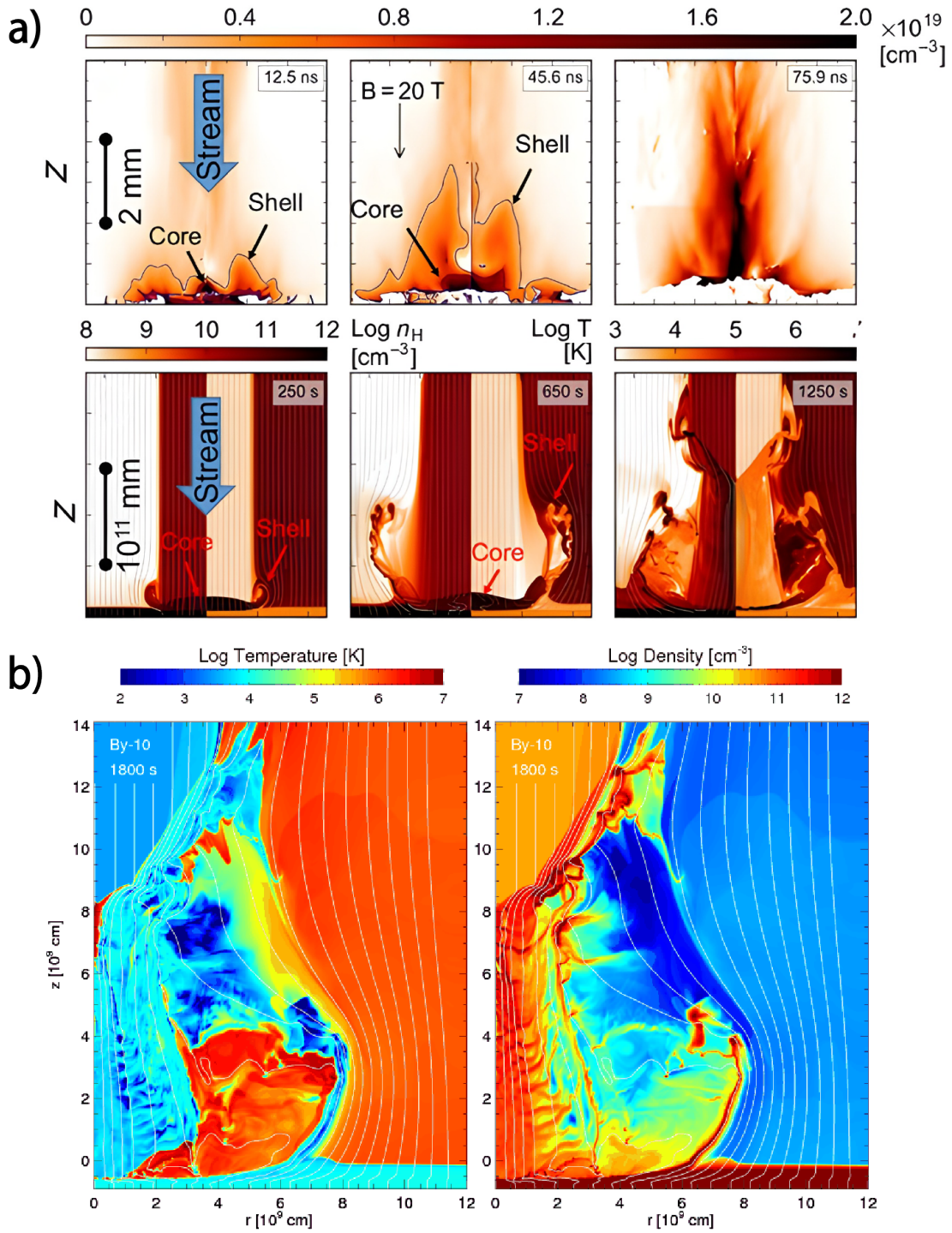


Figure 1.5: (a) Measured electron density maps of the magnetised laser-driven experiment presented in [Fig. 1.4(b)], showing the formation of a core region and a dense shell mediated by the magnetic field (top) along with results from numerical 2D astrophysical simulations (bottom). Reprinted with permission from [21] (b) Temperature and mass density distributions in (r, z) plane of the lateral post-shocked plasma interacting with a 10T magnetic field from 2D numerical radiative MHD simulations, reproduced with permission from Astronomy & Astrophysics, © ESO [26].

nature of these accretion shocks, where a dense shell region bounded by the magnetic field can form from the shocked materials escaping the core region [Fig. 1.5(a)]. These shells provide an extra source of X-ray absorption, thus, influencing the expected mass accretion of the stars.

This was also possible thanks to the extensive effort spent in numerical simulations of accretion shocks, which helped to complement and interpret the experimental observations. In particular, relevant numerical studies on accretion shocks are [26–28, 55–57]. Various magnetic field configurations have been explored in this context, including radiative effects on the stability and dynamics of accretion shocks in CTTs. Some of the key results from these studies produced by 2D radiative MHD models revealed that in cases of high plasma beta ($\beta = P_{th}/P_B \gg 1$) the post-shocked plasma can bend the magnetic field, creating a large component of the B perpendicular to the stream at the base of the accretion column, which limits the sinking of the slab into the chromosphere and perturbs the overstable oscillations of the shock [Fig. 1.5(b)]. In the case of a strong magnetic field ($\beta \ll 1$), the field configuration determines the position of the shock and its stand-off height. Moreover, non-uniform magnetic field strength can also produce oblique shocks at great distances from the chromosphere which can become unstable due to radiative cooling.

Although in the last decade, a few experiments successfully replicated the dynamics of accretion flows, most of the laser-driven experiments have focused on studying the X-ray emissivity and the accretion shock rate with very limited attempts to explore radiative cooling effects on the reverse shock structure, which still remains an open question. Replicating the geometry and the plasma conditions scalable to accretion shocks in young stellar systems is not trivial, and often radiative cooling phenomena occur at a similar timescale of the experiment, leading to challenges when it comes to studying the effect of radiative instabilities, such as thermal instability and thin-shell instability. The lack of laboratory analogous to radiative accretion shocks limits our current understanding of the absorption of the X-ray emitting plasma and the energy balance of the post-shock region, both important factors that are expected to possibly underestimate the mass accretion rates.

Moreover, long-lasting reverse shock systems in highly controllable environments are necessary to investigate the complex dynamic of accretion shocks. The accretion shock in the majority of astrophysical events has a much more complex structure than the classical quasi-one-dimensional

profile. The 3-dimensional nature of the system and the surrounding chaotic materials can well obscure the shock profile, appearing as a bright perturbed region. This is definitely the case in the study of

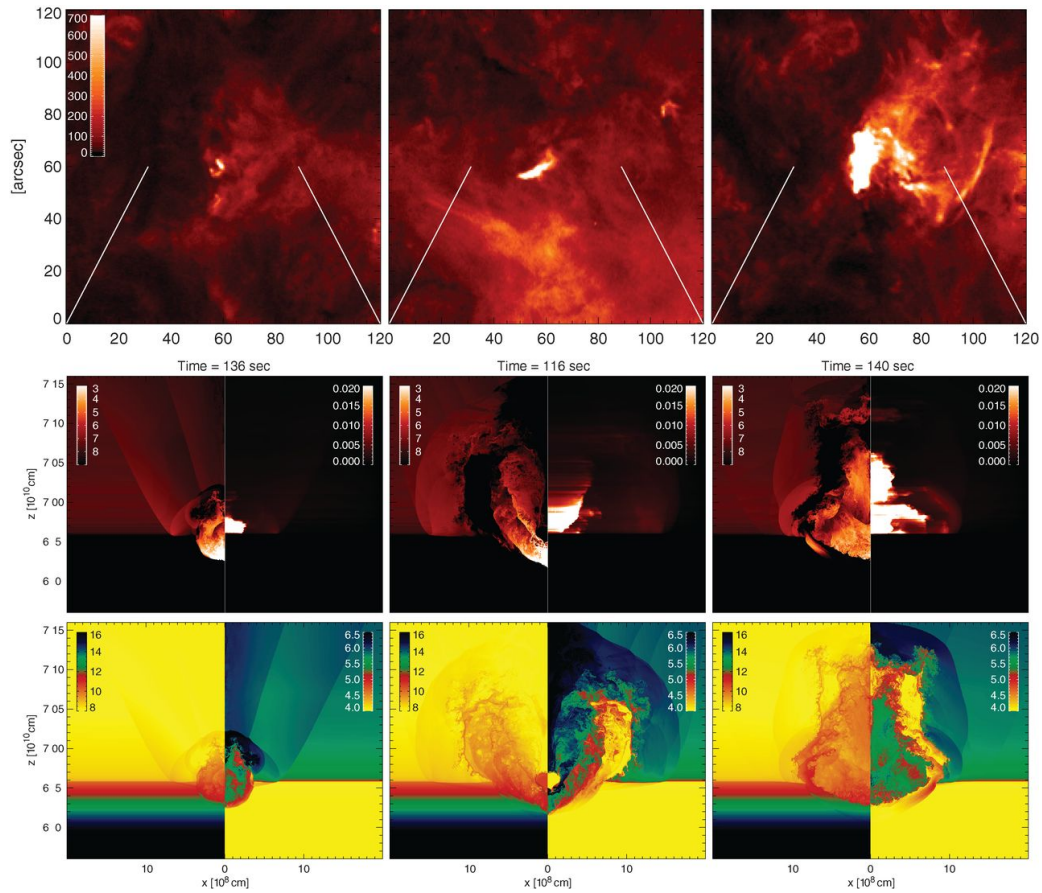


Figure 1.6: Results from the study on the impacts of dense fragments erupted after an M-class solar flare on 7 June 2011 in the SDO/AIA 171Å channel. The numerical simulations show the fragmented nature of the impact as a result of the interplay between hydrodynamics and radiation. Reprinted with permission from © The American Association for the Advancement of Science [58].

solar fragments [58]. Here, the study of falling plasma materials onto the chromosphere produced by a solar flare revealed the fragmented nature of the shocks, similar to what can happen in accretion shock in young stellar objects, Fig. 1.6.

In this thesis, an intense electric current pulse generated using a pulsed power generator drives the formation of a supersonic outflow from initially solid tungsten thin wires carrying an embedded magnetic field. A planar conductive obstacle is then placed perpendicular to the plasma flow to create a reverse shock system. This reverse shock accretes from the surface as more plasma material is accumulated in the post-shocked region, lasting for hundreds of nanoseconds and sustained by the continuous plasma injection. Combined with the open geometry and excellent diagnostic access in

these experiments, the long lifetime of the system allows a suite of non-perturbative, laser-based diagnostics to make detailed measurements of the key upstream plasma parameters. These experiments are built upon preliminary studies on MHD shocks where the use of low-Z elements for wires, such as aluminium, has shown the formation of an early stagnation shock near the surface and a later development of a stand-off shock purely mediated by the accumulation of magnetic flux ahead of the obstacle [59, 60]. The use of high-Z material for the wires allows access to dominant radiative regimes, forming reverse shock systems scalable to astrophysical accretion shocks.

In addition to exploring the structure of reverse shocks using a flat obstacle, a second set of experiments is presented to investigate accretion shocks formed from the collision of supersonic plasma flows in an ambient magnetic field. The counter-streaming plasma flows in this case are generated by X-ray ablation of solid silicon targets, allowing the study of completely free-boundary reverse shocks under various magnetic field configurations. A schematic illustration of the two experimental geometries is shown in Fig. 1.7. Despite the differences between the exploding wire system and the astrophysical case, the results of these experiments still provide valuable insights into the underlying physics and complex structure of accretion shocks with unprecedented details.

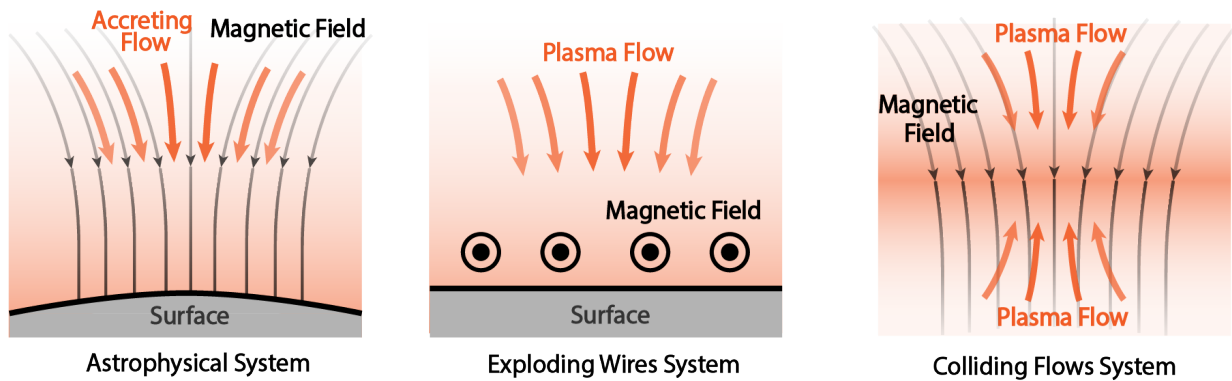


Figure 1.7: Schematic illustration comparing the astrophysical accretion shock system with the laboratory experiments presented in this thesis.

1.4 Outline of thesis

This thesis investigates the structure of reverse shocks and radiative cooling effects in experiments carried out at the MAGPIE facility at Imperial College London relevant to astrophysical accretion

shocks.

Chapter 1 presents a general introduction together with the motivations of this study. The basic theory of shocks is presented in §1.1 in the context of astrophysical accretion shocks commonly found in young stellar systems. This is followed by a discussion of the role of radiative cooling effects in §1.2, focussing on the two primary instabilities, thermal instability, and thin-shell instability.

Chapter 2 discusses a summary of the MAGPIE facility and the main diagnostics used throughout the investigations. Two different experimental platforms have been used for these experiments: an inverse wire array z-pinch and a newly developed X-ray driven platform using a standard wire array z-pinch, described in §2.3.1 and §2.3.2, respectively.

Chapter 3 introduces a novel refractive-based diagnostic, called imaging refractometer, including first a description of the optical configuration in §3.1 followed by a synthetic test case using a custom ray tracing tool in §3.2 and experimental validation by applying the diagnostic to a simple plasma object. Chapter 4 presents one of the two main experimental campaigns, looking at reverse shocks formed from the collision of super-magnetosonic plasma produced by an inverse wire array against a planar conductive obstacle. A summary of the experimental setup and diagnostics is provided in §4.1. Experimental results from reverse shock systems employing both aluminium and tungsten wires are discussed in §4.2 and §4.3.

Chapter 5 introduces the second experimental campaign on free-boundary reverse shocks from colliding counter-streaming plasma flows generated using the X-ray driven platform. Two ambient magnetic field configurations have been explored with results discussed in §5.1 and §5.2, respectively.

Chapter 6 presents the main conclusions and suggestions for future work.

1.5 Authors contributions

The experimental results presented in this thesis were conducted at the MAGPIE facility at Imperial College London, where I have been part of the research team since the start of my PhD. The team, consisting of Doctoral researchers and PhD students, is responsible for both the maintenance and development of the pulsed power machine and diagnostics. Each experimental campaign is the result of a team effort. I have been involved in the majority of the experiments conducted in the last four

years and I am a co-author on several publications [61–65]. Due to COVID-19 restrictions, I have been forced to re-balance the work towards data analysis of previous experiments and diagnostic development from March 2020 onwards. In that period, I contributed to the development of the synthetic shadowgraphy and imaging refractometer code presented in this work, with the support of Jack Hare and Aidan Crilly. I also carried out the analysis of the reverse shock experiments using exploding wire arrays conducted by Guy Burdiak, Jack Hare and Jack Halliday between 2016 and 2017. In 2021, I collected new experimental data using a re-designed wire array hardware, which is not included in this thesis, except for the optical self-emission images presented in Chapter 4. Although I was not directly involved in the experimental campaign led by Burdiak, I have actively contributed to the analysis, interpretation, and presentation of the data, including measurements of the novel imaging refractometer that led to my first publication as lead author [66]. The experiments at MAGPIE were also disrupted from August 2021 and almost the entire of 2022 due to major refurbishments of the pulsed power generator, involving transformer oil replacement, reparation of the vertical transmission line, as well as one of the "elbow" sections of the horizontal transmission line. In addition to these extraordinary refurbishments, I have also been part of the regular maintenance of Marx banks. The load hardware used in the experiments presented in Chapter 5 was designed by Jack Halliday. I supported the initial experiments using single ablating targets, which were published in [67], and collected the data for the free-boundary accretion shock experiments in Chapter 5. Additionally, I performed the HELIOS-CR [68] simulations presented in the discussion. Data analysis was carried out using several internal analysis codes developed mainly by George Swadling [69] and Jack Hare [70]. Selected results from this thesis work have been presented at the Annual Meeting of the APS Division of Plasma Physics 2020-2021-2022, IEEE Annual Conference Plasma Science 2022 (in person), IOP Annual Plasma Physics Conference 2023 and will be presented at the upcoming Annual Meeting of the APS Division of Plasma Physics 2023.

Chapter 2

The MAGPIE Pulsed Power Generator and Diagnostics

The MAGPIE (Mega Ampere Generator for Plasma Implosion Experiments) pulsed power generator is an experimental platform located at Imperial College London capable of delivering 1.4 MA of pulse current with 240 ns rise-time. This electric current is discharged in a load of a few centimetres, resulting in the creation of extreme physical states of matter. Initially designed for driving the implosion of a cryogenic hydrogen fibre in thermonuclear fusion research, MAGPIE is now used to generate hydrodynamic structures that resemble astrophysical objects [71–76]. The high impedance of the generator allows it to drive exploding and imploding targets with high inductance. This feature, combined with its versatility in reaching HED plasma conditions, offers significant advantages in investigating plasma flows under highly controlled environments.

Over the past few decades, a wide range of diagnostics has been developed to study ablated plasmas that are generated when the pulse current passes through solid loads. The extreme plasma states in MAGPIE make it difficult to probe with intrusive probes, such as Langmuir probes, therefore laser-based diagnostics are used for quantitative measurements of plasma temperature, density, flow velocity and magnetic field. While the author dedicated extensive time to maintaining the pulsed-power components, any facility improvements are beyond the scope of the current work, and therefore, only a brief description of the generator will be provided.

2.1 The MAGPIE Generator

The MAGPIE Generator comprises four Marx banks connected in parallel to the vertical transmission line (VTL), which terminates at the load vacuum chamber. The coupling between these elements is facilitated by four pulse-forming lines (PFLs), which compress the current pulse in time before charging the vertical transmission line (VTL). A schematic drawing of the facility is illustrated in Fig. 2.1, indicating the seven main components: four Marx banks consisting of six suspended capacitors connected by gas-filled switches, four PLFs, four trigatron switches, a single VTL, load vacuum chamber surrounded by an iron-curtain shielding and a data acquisition room.

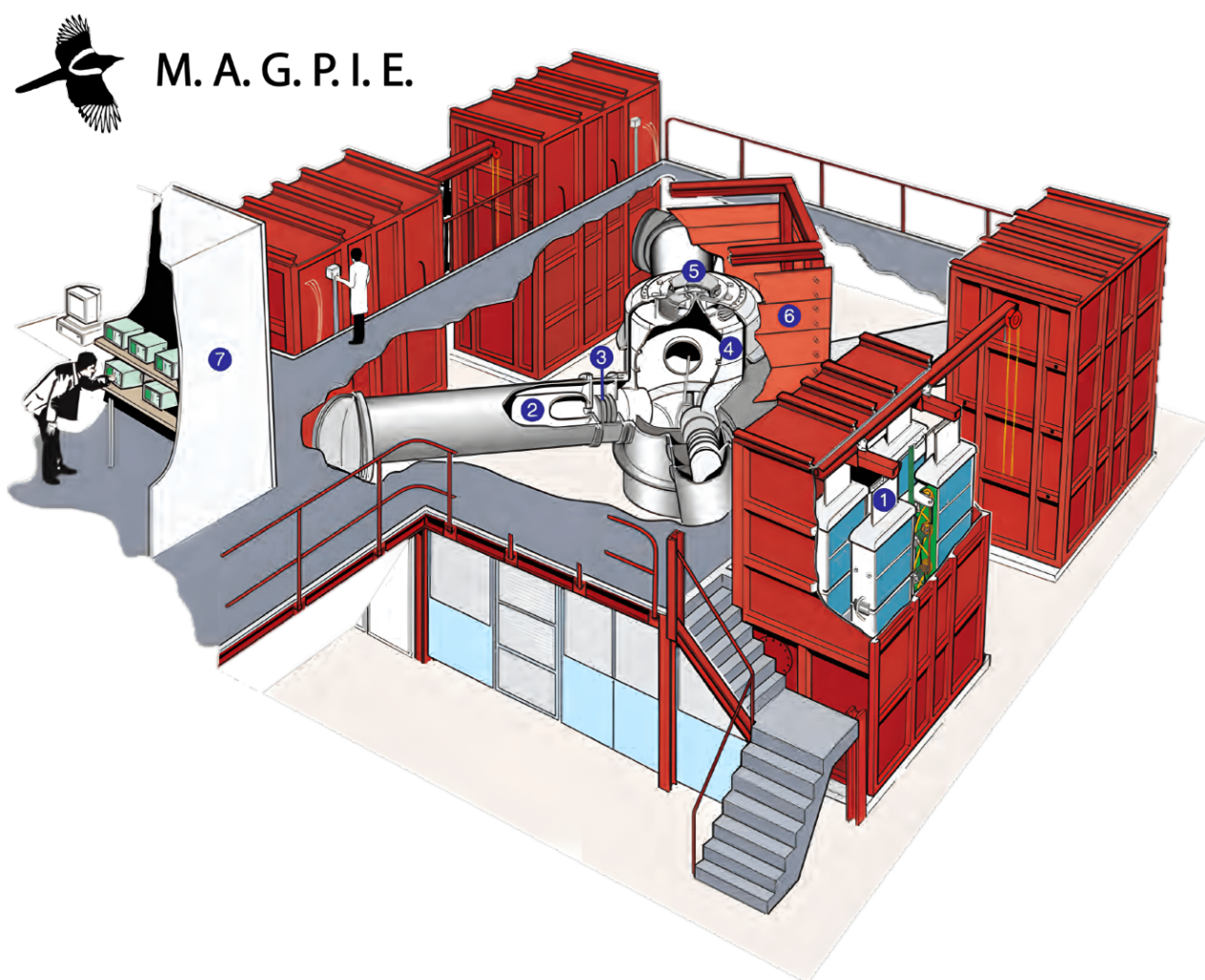


Figure 2.1: Drawing of the MAGPIE facility with cut sections to show the internal components. The labelled features are (1) Marx bank with an internal view of the suspended capacitors and switches. (2) Pulse forming line. (3) Trigatron switch. (4) Vertical Transmission line. (5) Vacuum load chamber. (6) Iron X-ray shielding. (7) Data acquisition room.

2.1.1 Marx Banks

Each Marx bank stores electrical energy in an array of $N = 24$ individual capacitors, each with a capacitance of $C = 1.3 \mu\text{F}$, connected in parallel. These capacitors are charged with alternating voltages of $\pm 65 \text{ kV}$ before being linked in series by the breakdown of twelve SF_6 gas-filled switches. In this configuration, the capacitor voltages are stacked to achieve a high voltage signal of 1.5 MV . The total energy stored in all 4 Marx banks is 264 kJ . An external trigger Marx (300 kV , 700 J) drives the erection phase, producing an electric discharge on the first five spark-gap stages, leading to the charge of the pulse forming lines.

2.1.2 Pulse Forming Lines

The pulse-forming lines are large coaxial transmission lines filled with deionised water which acts as dielectric material increasing the capacitance for the storage of the energy from the Marx banks. Moreover, the high resistivity ($10^4 \Omega\text{m}$) of this medium prevents undesired breakdowns. This intermediate energy storage is isolated from the VTL via further spark gaps known as trigatrons. The PFLs are charged throughout $\simeq 1 \mu\text{s}$ and discharged upon the firing of the trigatrons at the opposite side of the line. The discharge phase towards the VTL happens over 240 ns , effectively compressing the available energy in time. The inductance and capacitance of the PFLs set the amplitude of the MAGPIE current pulse and thus, the overall rise time. A linear current rise time is not achieved on MAGPIE and the real current pulse delivered to the load is better approximated as follows:

$$I_{load}(t) = I_0 \sin^2\left(\frac{\pi t}{2t_{rise}}\right) \quad (2.1)$$

where the I_0 is the peak current and the $t_{rise} = 240 \text{ ns}$ is the rise time of the pulse.

2.1.3 Vertical Transmission Line, MITL and Vacuum Load Chamber

When the trigatrons are closed, a single current pulse is formed passing through the vertical transmission line. The VTL is a vertical coaxial transmission line with an outer diameter of 2 m filled

up with deionised water as internal dielectric material, resulting in an impedance of 1.25Ω . The main function of the VTL is to ensure that a uniform signal wave-front reaches the MITL. Here, the current pulse is compressed in space rather than in time, focusing the electron current from the 2 m VTL diameter into a few centimetres of the load size. An illustration of the MITL is represented in Fig. 2.2.

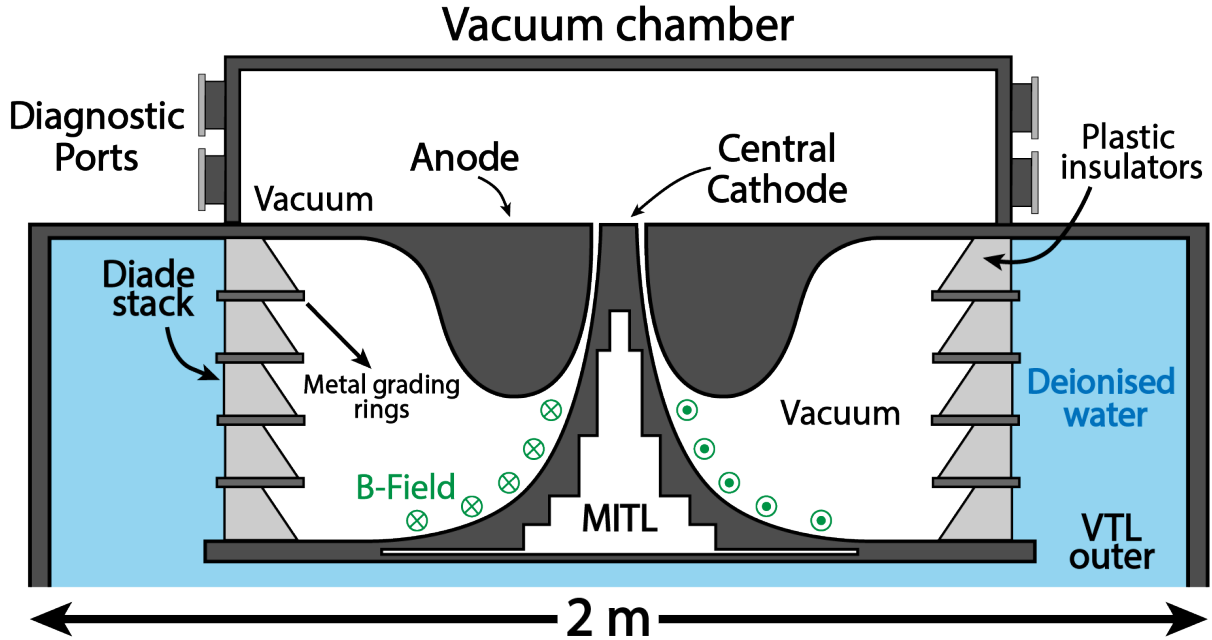


Figure 2.2: Schematic representation of the vacuum power feed, consisting of the MITL and the vacuum load chamber. Deionised water fills the gap between the VTL outer and inner conductors and also the hollow region below the MITL. The diode stack separates the water from the vacuum which fills the internal volume of the MITL and the load chamber. The structure of the diode stack and the MITL determines the shape of the electric field and prevents electrical breakdown.

Although the current compression is performed in a vacuum, the presence of a large potential between the two coaxial conductors is normally problematic. This is overcome due to the formation of a large azimuthal magnetic field generated when the current passes through the MITL. Under this condition, electrons that are liberated from the surface by a large electric field cannot reach the anode-cathode gap due to the magnetic field which pushes them back to the same surface by the drift $\mathbf{v} \times \mathbf{B}$. This magnetic insulation becomes stronger as the diameter of the inner conductor decreases, $B \propto 1/r$, until the current peak is reached. The interface water/vacuum is given by an alternating structure of metal and plastic rings forming the diode stack. This arrangement allows the shaping of the electric field in the inner region such that electrical breakdowns do not occur.

The final stage of the MAGPIE generator is a vacuum load chamber where the current pulse is

delivered to the target by means of a coaxial anode-cathode platform. This arrangement is usually loaded with a solid material such as thin wires or foils which are converted into highly dense plasmas by Ohmic heating as the current pulse flows through the load from the anode to the cathode. The two platforms used in this work are described in § 2.3.1 and 2.3.2 which are mounted on top of the anode and cathode structures in Fig. 2.2.

2.2 Diagnostics

The MAGPIE load vacuum chamber is characterised by multiple vacuum windows distributed on both side surfaces and lid. This flexible design accommodates a wide variety of diagnostics capable of probing from several direction angles. Among standard voltage and current probes useful to monitor the performance of MAGPIE during a shot, highly developed laser-based diagnostics provide quantitative insight into the plasma state. Two-colour laser interferometry and Faraday rotation imaging allow measuring plasma electron density and magnetic field respectively, while electron temperature and flow velocity are derived from an optical Thomson Scattering system. For qualitative images, refractivity-based diagnostics, such as shadowgraphy and schlieren imaging, are used to capture shocks and density structures in plasmas. Although those diagnostics can give us highly accurate measurements of plasma properties within the time scale of a MAGPIE experiment, retrieving quantitative data from strong density gradient perturbations is still a non-trivial problem and further improvements of the current optical systems must be accomplished. In the next chapter, a novel refractive-based laser diagnostic, named imaging refractometer, is introduced as a new technique for quantitative measurements of electron density gradients in highly perturbed plasmas. This section is organised as follows: first, a short description of machine diagnostics and optical self-emission imaging is presented. Second, the main refractive-based laser diagnostics are introduced, such as imaging interferometry, shadowgraphy and schlieren. Finally, the optical Thomson scattering diagnostic is discussed in detail, completing the overview of the measurement techniques used in this work.

2.2.1 Load Current Diagnostics

The load current delivered to the experiment fielded on a pulsed-power generator represents an important quantity to diagnose, indicating the amount of power discharged through the platform during the entire duration of the experiments. In order to characterise accurately the load current, experiments on MAGPIE use current probes named Rogowski coils, capturing the inductive pickup of the time-varying magnetic field generated by the current pulse. Fig. 2.3 shows a cross-section schematic of a Rogowski probe around a conducting post. The magnetic field induced by the current flowing through the central conductor, $I(t)$, is linearly proportional to the current as:

$$B = \frac{\mu_0 I}{2\pi r} \quad (2.2)$$

The change in magnetic flux due to the change of magnetic field over the area of the groove generates a voltage signal which is detected by a coaxial cable and then recorded using an oscilloscope. Since this voltage is given by $V = -d\Phi/dt$, the signal is therefore proportional to the derivative of the current, dI/dt as follows:

$$V(t) = \frac{\mu_0 h \ln(r_2/r_1)}{2\pi} \frac{dI(t)}{dt} \quad (2.3)$$

Where r_1 , r_2 and h are the dimensions of the probe as shown in Fig. 2.3. For reference, the standard Rogowski probes used in the experiments have the following dimensions $r_1 = 13$ mm, $r_2 = 9$ mm and $h = 3.5$ mm. Integrating the voltage signal over time leads to the measurement of the load current pulse. This is usually performed numerically rather than using a passive electrical integrator. It is important to note that the Rogowski probes are sensitive to the total current so they cannot provide information on the current distribution. At MAGPIE two Rogowski probes are normally applied to two opposite return posts in each experiment. The use of two probes with opposite polarity enables the removal of any unwanted capacitively coupled voltages affecting the signal dI/dt . In this case, the combined voltage will be:

$$\bar{V}(t) = \frac{V_+(t) - V_-(t)}{2} \quad (2.4)$$

with $V_+(t)$, $V_-(t)$ the voltage signals of the two probes. The contribution of the capacitive pickup affecting the voltages is the same for both probes with opposite signs so that when combining the two

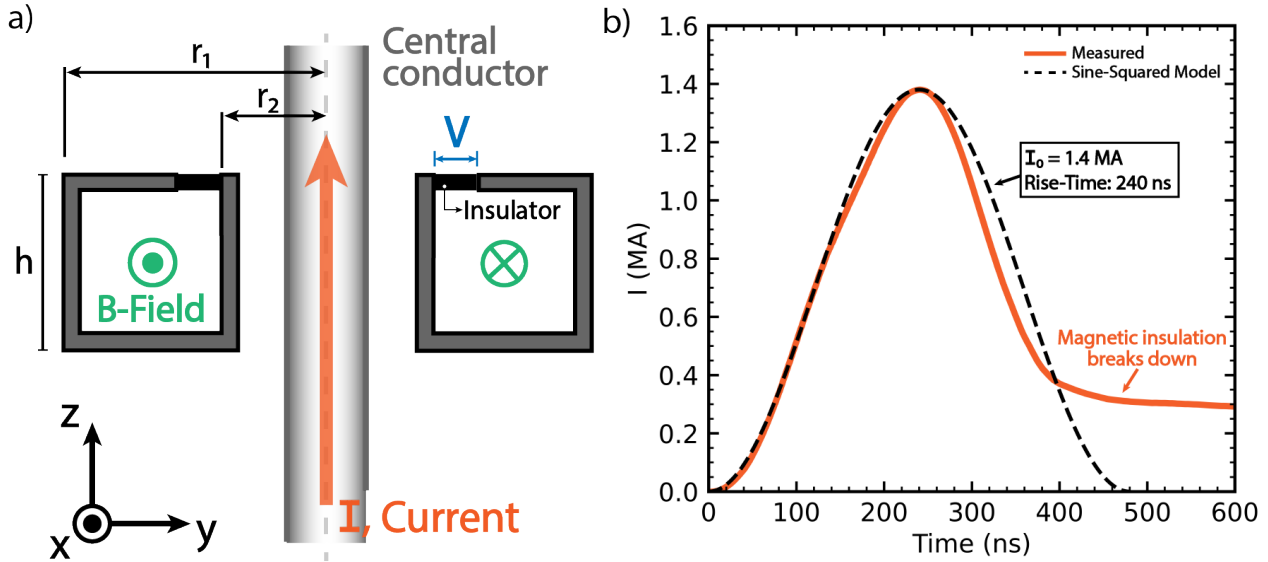


Figure 2.3: a) Schematic representation of a Rogowski probe around a central cathode post. The probe surrounds the carrying current conductor. b) An example of the electric current profile measured during an experiment (orange), following a sine-squared model of Rise-Time 240 ns and a peak current $I_0 = 1.4$ MA.

signals, this capacitive effect cancels out. An example of the load current profile measured from one of the experiments is shown in Fig. 2.3. The starting time of the experiment in each shot is defined by the first time the current is detected by the Rogowski probes. The load current is well approximated by the sine-squared profile eq. (2.1) until the MITL breaks down. In addition to the amount of energy expected in the load, the Rogowski probe signals in combination with the other machine diagnostics provide an assessment of the reproducibility of the shot. Depending on the sensitivity of the experiment to the current profile, shots in which the current profile is not ideal may be discarded.

2.2.2 Optical Self-Emission Fast Framing Camera

Optical self-emission diagnostics are useful techniques that image the plasma emission in the optical wavelength. The primary optical self-emission diagnostic used at MAGPIE is a fast-framing camera able to capture a sequence of 12 images with a 5 ns exposure time. This camera records the optical broadband emission of the plasma over a wide temporal range set by the interframe between two consecutive images. In the present work, the interframe is typically between 25 - 40 ns corresponding to a 300 - 480 ns period. This time-resolved capability makes this diagnostic excellent for observing the

evolution of structures within the time scale of the experiments. The camera uses the same line-of-sight of the refractive-based laser diagnostics so that it is easier to relate features across the camera field of view between the optical self-emission images and the laser probe diagnostics, specifically with the quantitative imaging interferometry. The optical arrangement consists of a 10:1 telescope positioned immediately out of the vacuum chamber composed of two lenses with focal lengths, $f_1 = 750$ mm and $f_2 = 75$ mm. A notch filter that reflects wavelengths between 500 nm and 560 nm is placed just in front of the camera, removing the probe laser from the self-emission images. In general, plasma self-emission can depend on various plasma parameters, like temperature, density and charge state, however for the plasma driven at MAGPIE, this primarily originated from line emission and so is proportional to the square of the plasma density.

2.2.3 Refractive-based Laser Diagnostics

At MAGPIE, various laser probing techniques are employed to image the plasma and obtain both qualitative and quantitative information about its state. One of the benefits of using lasers instead of physical probes is that they are often non-invasive, meaning they do not alter the system by introducing any external disturbances. As already mentioned, MAGPIE has several laser-based diagnostics that can measure plasma properties such as density, temperature, velocity, and magnetic field. These diagnostics are pulsed and triggered once per shot at a specific time. In this section, we focus on laser techniques that depend on the refractive index of the plasma, specifically shadowgraphy, schlieren and imaging interferometry while optical Thomson scattering diagnostic will be discussed separately.

The refractive index, N , is a physical property of the material that describes how it affects the propagation of electromagnetic (EM) waves. It is defined as the ratio between the speed of light in vacuum, c , and the phase speed of the EM waves in the material, v_{ph} :

$$N = \frac{c}{v_{ph}} \quad (2.5)$$

$$v_{ph} = \frac{\omega}{k} \quad (2.6)$$

where k is the wavenumber and ω is the frequency of the laser beam. In plasmas, the phase speed at a

specific wavelength is greater than c , while in the neutral medium is lower than c . In this case, the phase velocity, v_{ph} , is defined from the dispersion relation (ω - k relationship) for EM waves propagating through a plasma medium. Assuming an isotropic plasma with no magnetic field ($B = 0$), the dispersion relation can be derived as a particular solution of Ampere's and Faraday's Laws:

$$\frac{\partial \mathbf{B}}{\partial t} = -\nabla \times \mathbf{E} \quad (2.7)$$

$$\nabla \times \mathbf{B} = \mu_0 \mathbf{j} + \frac{1}{c^2} \frac{\partial^2 \mathbf{E}}{\partial t^2} \quad (2.8)$$

By taking the time derivative of (2.8) and substituting in the curl of (2.7), the general wave equation for a given electric field, E , is therefore expressed as:

$$\nabla^2 \mathbf{E} - \nabla(\nabla \cdot \mathbf{E}) = \mu_0 \frac{\partial \mathbf{j}}{\partial t} + \frac{1}{c^2} \frac{\partial^2 \mathbf{E}}{\partial t^2} \quad (2.9)$$

Assuming a planar wave, $\mathbf{E}$, travelling in $\hat{z}$ direction with the electric field in the $\hat{x}$ direction:

$$\mathbf{E} = E_x e^{i(kz - \omega t)} \hat{\mathbf{x}} \quad (2.10)$$

and carrying on with the differentiation gives the relation:

$$-k^2 \mathbf{E} = \mu_0 \frac{\partial \mathbf{j}}{\partial t} - \frac{\omega^2}{c^2} \mathbf{E} \quad (2.11)$$

Now, the current is only dependent on the electron motion due to the electric field if ion and thermal motion are neglected, hence:

$$j_x = -n_e e v_x \quad (2.12)$$

with n_e is the plasma electron density, e is the electron charge and v_x is the electron velocity. Substituting j into (2.11), we obtain the dispersion relation for a plasma, remembering that $\epsilon_0 \mu_0 = 1/c^2$ and the plasma frequency $\omega_p^2 = n_e e^2 (\epsilon_0 m_e)^{-1}$:

$$\omega^2 = \omega_p^2 + c^2 k^2 \quad (2.13)$$

According to (2.13) the definition of phase velocity in a plasma can be written as:

$$v_{ph} = \frac{\omega}{k} = \frac{c}{\sqrt{1 - \omega_p^2/\omega^2}} \quad (2.14)$$

and thus, the plasma refractive index:

$$N = \sqrt{1 - \omega_p^2/\omega^2} = \sqrt{1 - n_e/n_{cr}} < 1 \quad (2.15)$$

In many laser diagnostic applications, the derived refractive index, (2.15) is usually expressed in terms of the critical density $n_{cr} = \omega^2 \epsilon_0 m_e / e^2 \simeq 1.21 \times 10^{21} \lambda_0^{-2}$ where λ is the free space frequency in μm of the EM beam. This critical density represents the limit on whether or not the laser beam can propagate through the plasma. In both laser wavelengths used at MAGPIE, 532 nm and 355 nm, this critical electron density is very high and comparable density conditions are only found in the stagnation column of a wire array Z-pinch implosion [69, 77] or very close to the ablating wire cores. The experiments presented in this work are unlikely to exhibit such high densities, hence, the laser diagnostic capability is instead determined by the strength of electron density gradients that refract the laser light out of the collection optics.

In the presence of magnetic fields, the refractive index depends upon the orientation of the magnetic field lines and the direction of the EM wave propagation. Magnetic fields affect the motion of the electrons, altering the total current density $\mathbf{J}$ through changes in the plasma conductivity. This results in a quadratic equation for N^2 , named the Appleton-Hartree formula, which is dependent upon both ω_p^2/ω^2 and Ω/ω , where $\Omega = eB_0/m_e$ denotes the electron cyclotron frequency proportional to the magnetic field, B_0 . Plasmas generated at MAGPIE are usually magnetised, however, the magnetic field is not strong enough to have relevant effects on the refractive index, since $\Omega/\omega \sim 3 \times 10^{-4} \ll 1$. A more detailed derivation of the refractive index including magnetic field effects can be found in [78].

Shadowgraphy & schlieren imaging

A wide range of optical diagnostics is based on the refractivity principle. Shadowgraphy and schlieren techniques have been developed since the early studies of experimental hydrodynamics,

starting from flash lamps and photo-sensitive films to collimated laser beams and high-resolution cameras [79]. These refractivity-based techniques are used in plasma physics to image very non-uniform electron density regions (hence non-uniform refractive index). The physical principle behind both shadowgraphy and schlieren imaging is that light rays are deflected by refractive index gradients normal to their propagation direction. For a single path through the plasma, the integral of each infinitesimal element dz with its own index of refraction $N(z)$ represents the total deflection angle of the ray ϕ (radians) due to lateral gradients of the plasma electron density:

$$\phi = \int \nabla N dz = \frac{1}{2} \int \frac{\nabla n_e}{n_{cr}} dz \quad (2.16)$$

This redistribution of light rays results in a light intensity modulation at the image plane of the camera system.

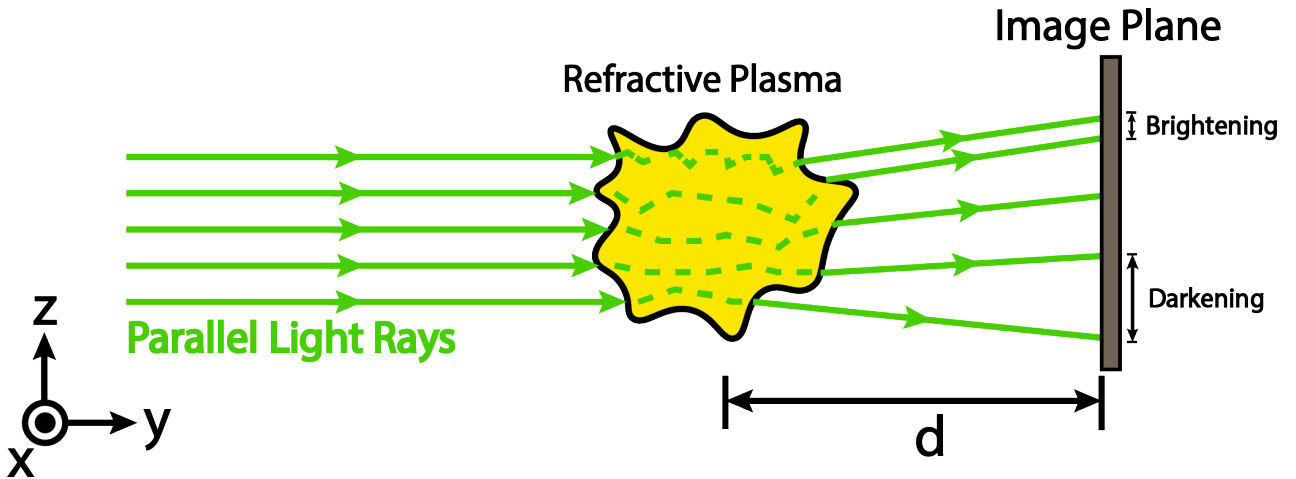


Figure 2.4: Schematic illustration of shadowgraphy method. Light of the laser probe beam is redistributed at the image plane due to the passage through a refractive plasma. This results in an image with dark and bright regions depending on the second spatial derivative of the plasma electron density.

Fig. 2.4 represents a schematic illustration of the shadowgraphy method applied to a non-uniform plasma. The “shadow” of the plasma formed at distance d is sensitive to the second spatial derivative of the refractive index expressed by the following relation:

$$\frac{\Delta I}{I} = L \left[\frac{d^2}{dx^2} + \frac{d^2}{dy^2} \right] \int N dz \quad (2.17)$$

Where $\frac{\Delta I}{I}$ is the light intensity variation at the detector and L is the spatial length of the plasma. For regions where the electron density gradient is increasing (the second derivative is positive), the plasma behaves as a diverging lens and the image appears darker, while decreasing gradients cause the rays to converge, appearing brighter in the image.

Unlike shadowgraphy, schlieren imaging is sensitive to the first spatial derivative of the refractive index. A typical configuration is shown in Fig. 2.5 which consists of a typical shadowgraphy system with a circular disk placed at the focus point. This beam stop prevents all the undeflected laser lights from reaching the camera. The presence of a non-uniform plasma electron density distribution deflects light rays due to lateral refractive index gradients according to the equation (2.17). The minimum gradient (cut-off) is set by the size of the beam stop while symmetry allows gradients in all directions to be imaged. Although schlieren images contain the location of density gradients sampled by the probing beam, almost no information can be retrieved about their magnitude. Both shadowgraphy and

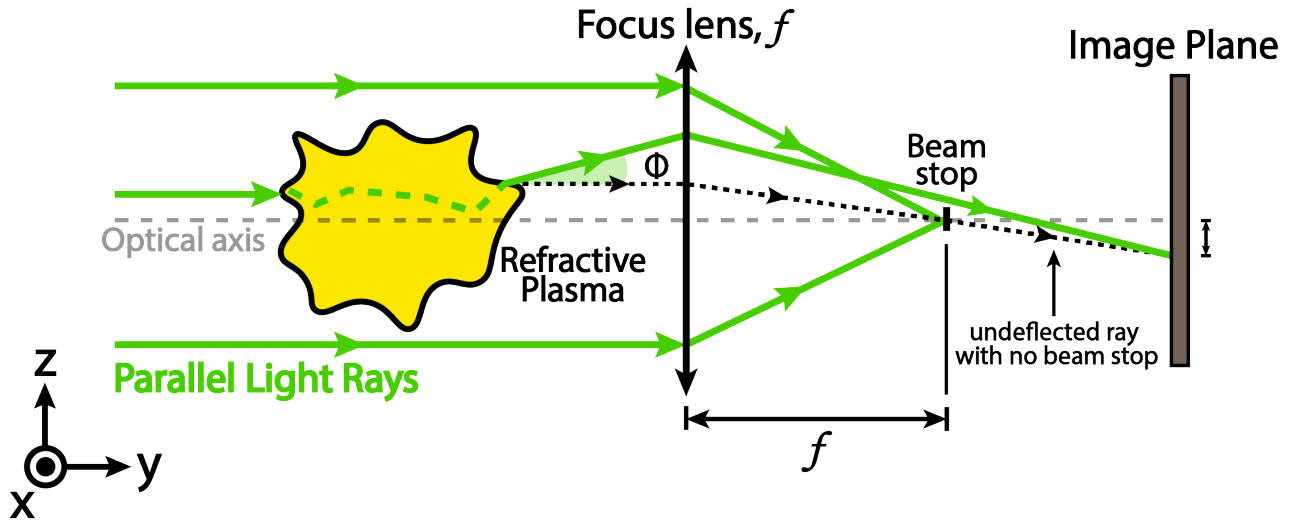


Figure 2.5: Schematic illustration of a schlieren imaging system with a single focus lens, f . Parallel light rays are refracted by the plasma and then collected by a spherical lens. A small object of known size is placed at a distance f from the lens so that only rays refracted by angles greater than ϕ can reach the detector. The image plane is formed at twice the focal length of the focus lens.

schlieren imaging have been widely used in both conventional hydrodynamics and plasma physics to reveal structures in flows due to high-density perturbations. For this application, the analysis of turbulence consists of a post-processing Fourier transform of the “shadow” images, leading to the intensity spectrum. This is linked to a spectrum of density fluctuations, so a power-law fit is performed and compared to the theoretical predictions. This method has one significant drawback, the dynamic

range is restrained by the detector pixel size, limiting the smallest resolvable length scale. In this condition, it is difficult to fit a power-law spectrum since the size of the driving scale and the smallest scale are often comparable (single order of magnitude). Moreover, large density gradients can result in ray crossings which leads to caustics, preventing the reconstruction of a unique profile of the density fluctuations.

All these limitations narrow down the range of cases where these techniques are used for quantitative investigations. Recent attention has been focused on possible extensions of the latter methods in order to extract meaningful plasma density fluctuations data [80–82], however, it is still an open challenge. In this work, the shadowgraphy technique is used to qualitatively image plasma structures and as a complementary diagnostic in support of the novel imaging refractometer diagnostic which will be discussed in detail in Chapter (3).

Imaging Interferometry

Imaging interferometry is a technique that uses the interference of light waves to measure the line-integrated electron density of a plasma. When a laser beam propagates through a plasma, phase changes in the beam can occur due to changes in electron density on the planes along the probing direction. The basic principle is to split a laser beam into two parts: a probe beam that passes through the plasma and a reference beam that does not. The two beams are then recombined to form an interference pattern which is recorded by a CCD camera. The presence of plasma along the probing path induces laser phase shifts relative to the reference beam. This phase difference between the probe and reference beams depends solely on the refractive index of the plasma, which depends on the plasma electron density. In its simple form, the imaging interferometry setup used at MAGPIE is shown in Fig. 2.6 which uses a Mach-Zender interferometer, as described in [78]. This configuration allows the recording of a 2D phase map of the interference pattern formed by the recombination of the reference and the probe beams. The incident laser beam is separated by a 50:50 beamsplitter forming a reference beam and a probe beam, orthogonal to each other. The probe beam is brought onto the vacuum chamber and through the plasma by a sequence of mirrors until the recombination beamsplitter which then interfere with the reference beam. Since the two laser paths are not equal, the

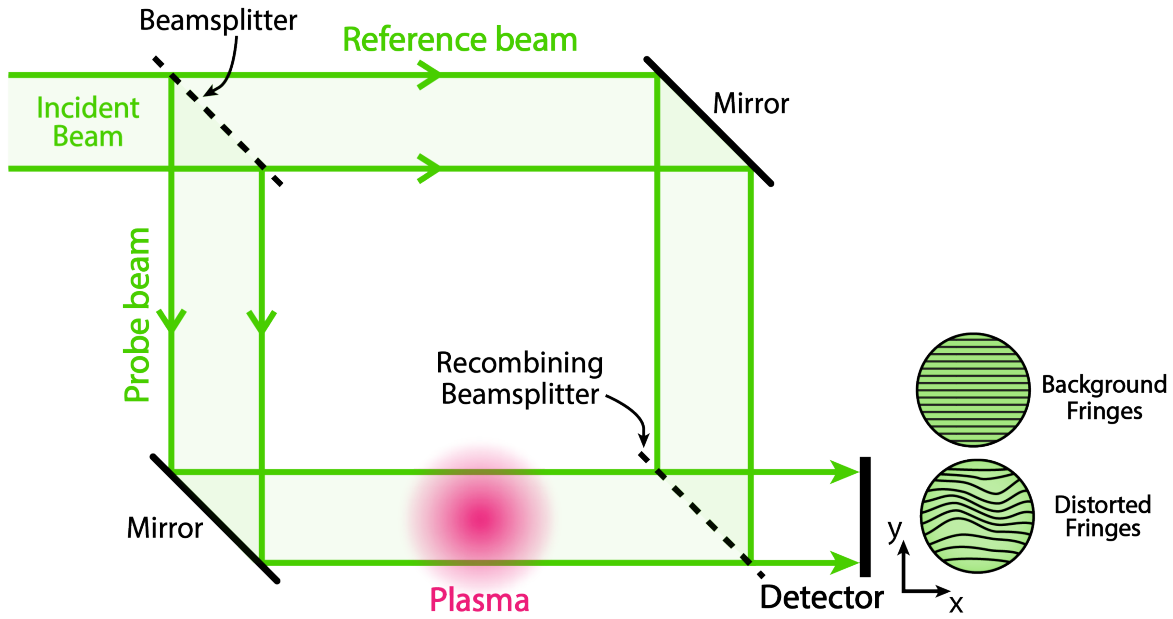


Figure 2.6: Schematic illustration of a Mach-Zender interferometry used to measure plasma electron density. The system consists of a single beam which is split into two orthogonal beams, a reference beam and a probe beam passing through the plasma. The reference and the probe are then recombined, forming an interference pattern which is finally imaged at the detector.

initial phase distribution can be determined from a pre-shot background image in the absence of any plasmas. The presence of an initial phase shift results in a constructive and destructive interference pattern represented by dark and bright fringes at the image plane, forming the final interferogram. Each fringe in the background interferogram indicates a constant phase and increases by a multiple of 2π for consecutive fringes, so that the spatial phase distribution, $\phi(x, y)$, is determined.

During the experiment, the phase distribution ϕ , changes as the laser probe interacts with the plasma producing a displacement of the interferogram fringes relative to their original background positions. The distorted pattern is referred to as the shot interferogram, exhibiting fringe displacements which are proportional to the plasma electron density. This phase shift can be quantified by comparing directly the background and the shot interferograms.

The phase, ϕ_p , accumulated by a monochromatic beam propagating through an unmagnetised plasma is given by the dispersion relation derived earlier (eq. 2.13), as follows:

$$\phi_p = \int \mathbf{k} \cdot d\mathbf{l} = \int \sqrt{\frac{\omega^2 - \omega_{pe}^2}{c^2}} dz \quad (2.18)$$

In the experiments presented here, the electron plasma frequency, ω_{pe} , is small compared to the frequency of the laser probe, $\omega \gg \omega_{pe}$ so that the expression (2.18) can be Taylor expanded obtaining:

$$\phi_p = \int \sqrt{\frac{\omega^2 - \omega_{pe}^2}{c^2}} dz \approx \int \sqrt{\frac{\omega^2}{c^2} \left(1 - \frac{\omega_{pe}^2}{\omega^2}\right)} dz \quad (2.19)$$

Knowing that the accumulated phase, ϕ , for the background interferogram, is simply given by integrating the phase velocity, v_p , along the probe direction, the total phase difference is calculated as:

$$\Delta\phi = \phi_p - \int \frac{\omega}{c} dz = - \int \frac{\omega_{pe}^2}{2c\omega} dz = \boxed{-\frac{e^2\lambda}{c^2m_e} \int n_e dz} \quad (2.20)$$

Where λ is the laser wavelength, and m_e and e are respectively the mass and charge of the electrons. This derived relation demonstrates the capability to estimate the line-integrated electron density from measurements of fringe displacements of the interferograms.

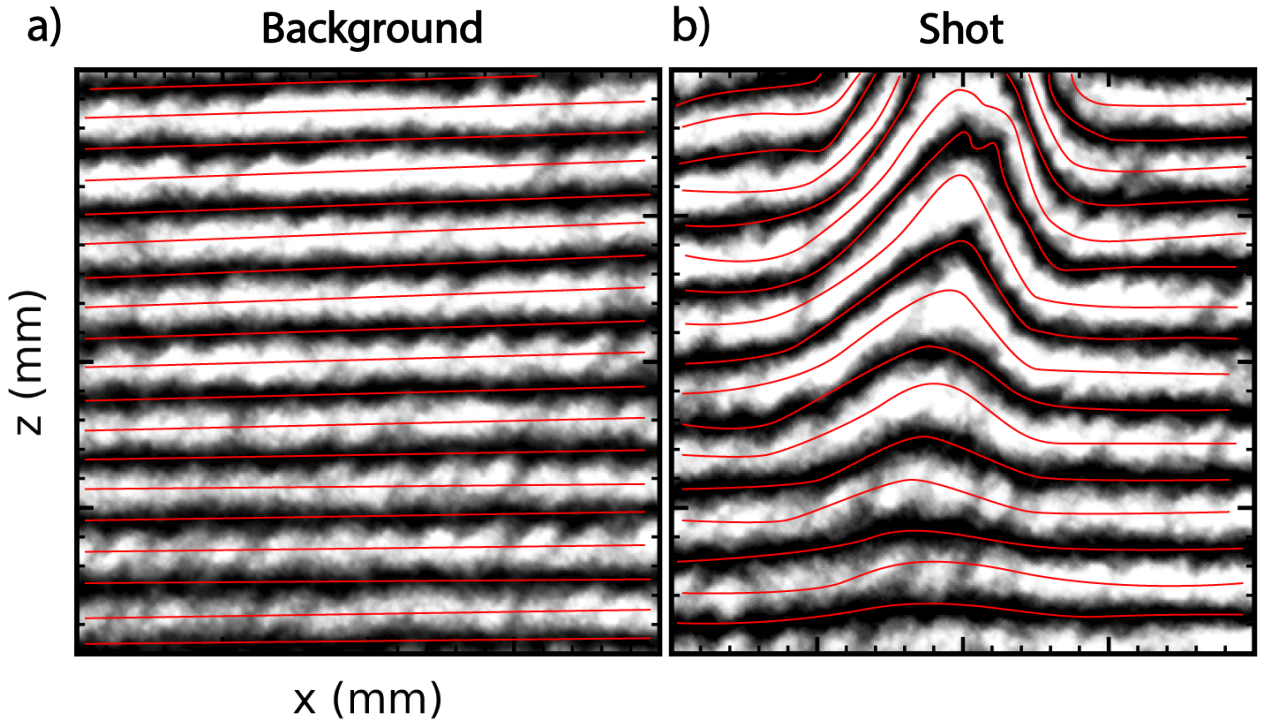


Figure 2.7: (a) Background interferogram taken before the shot. (b) Shot interferogram taken during the experiment in the presence of plasma. The fringe displacements are calculated by eq. (2.20)

An example of experimental interferograms is shown in Fig. 2.7. The shot interferogram shows the displacement of fringes along the z direction due to the presence of a plasma column at the image centre, while no fringe shift is recorded outside the plasma region, remaining unchanged relative to

the background, Fig. 2.7 (a). The process of measuring the fringe displacement from interferograms involves three main steps: first, both bright and dark fringes are traced manually using an image editor (e.g. Photoshop). Automated tracing is usually possible for the background images, however, the shot image presents often steep inflexions and complex features making the automation impractical. The results of the tracing are binary images of the background and the data interferogram with a value of 1 at the location of each fringe and 0 everywhere else. Second, the traced fringes are labelled such that the fringes in the shot interferogram are matched up with their corresponding fringes in the background and then interpolated to produce the respective 2D phase maps. Finally, the shot and background phase maps are subtracted according to eq. (2.20) to compute the 2D fringe displacement map, $\Delta\phi(x,y)$. The labelling method and the phase shift calculation were originally implemented in a *MATLAB* code called MAGIC [69, 72], and recently updated to MAGIC2 for *Python* [70]. MAGIC2 interpolates the phase relation of the traced images using a Delaunay triangulation algorithm to calculate the phase map $\phi = \phi(x, y)$ [72]. In order to calculate the absolute fringe shift, a region of zero-fringe displacement must be present in the shot interferogram, corresponding to zero electron density. In case this is not available, the relative fringe shift can still be measured, hence the relative electron density with respect to the minimum fringe shift. The unknown error between the assigned zero and the true zero is referred to as the "zeroing error". The angle and the separation of the fringes depend on the angle formed between the probe and reference beams at the recombination beam splitter. This angle lies within a specific plane and is chosen in advance before the experiment. The spatial sensitivity of the imaging interferometry to measure electron density changes is defined by the spacing between consecutive fringes. Measuring fringe displacement smaller than one width shift is challenging, therefore the spacing must be sufficiently tight to ensure good spatial sensitivity. Although space widths comparable to the pixel size of the CCD are theoretically achievable, in practice too closely spaced fringes may become difficult to trace when distorted in the shot image, thus limiting the maximum phase measurement of the interferogram. To further optimize sensitivity at the density gradient of interest, the fringes are typically oriented to be approximately perpendicular to the expected direction of density gradients in the experiment. The uncertainty of the measurement is usually given by the thickness of the fringes, which limits identifying the location of destructive interference used to calculate the phase shift distribution. This is explained in [83] where a 2π displacement of the

dark fringes in the plasma interferogram results in a line integrated density of $\simeq 10^{17} \text{ cm}^{-2}$. Visually tracking the minima by hand is accurate to about 1/5 fringe separation, thus the diagnostic uncertainty is of the order of the lower detectability limit of $\simeq 10^{16} \text{ cm}^{-3}$.

At MAGPIE, three Mach-Zehnder interferometers are commonly used. One of them is integrated into the Faraday rotation imaging diagnostic which employs a long-pulse arm of the Nd:glass CERBERUS laser ($\lambda_0 = 1053 \text{ nm}$, 100 mJ, 1 ns). The remaining two interferometers use an EKSPLA laser (Nd:YAG, 1064 nm, 60 mJ, 150 ps) as the laser pulse source. Frequency doubling and tripling crystals generate the second (green, 532 nm) and third (UV, 355 nm) harmonics which are then split and used in the two-colour interferometry systems for the two orthogonal directions, [70]. The beam diameter is approximately 30 mm and is split into a reference and a probe beam by a 50:50 beamsplitter. The probe beam is then brought into the chamber by a sequence of optical elements, including collimating telescopes and mirror, which passes through the load from both "end-on" (parallel to the axis of the wire array) and "side-on" (perpendicular to the axis of the wire array) directions. Two achromatic lenses with focal lengths, $f_1 = 750 \text{ mm}$ and $f_2 = 75 \text{ mm}$ relay an image of the laser beam at the centre of the plasma object to a diagnostics bench where the beam is then recombined with the reference beam. A two-lens imaging system composed of $f_1 = 500 \text{ mm}$ and $f_2 = 300 \text{ mm}$ lenses demagnifies the recombined beam to the size of a DSLR CCD camera. Commercial DSLRs (Canon 500D) represent the ideal cameras for these systems. They are robust to EMPs produced by the generator, have large CCDs and can be easily modified to be sensitive to IR and UV laser light. The interferograms are recorded by keeping the camera shutter open from the moment the generator is triggered until the end of the experiment (camera shutter opening time = 1.3 s). The time resolution of the interferograms is therefore set by the duration of the laser pulse, 150 ps.

2.2.4 Optical Thomson Scattering

Until now, changes in the properties of electromagnetic waves propagating through a plasma medium, such as phase shift, have been considered. However, the interaction between laser light and charged particles can result in scattering effects which are strongly coupled with the plasma state. The spectral properties of the scattered light carry information on the plasma state from which it has

is elastic which results in a very little energy transfer between the photons and the electrons. This has two consequences: firstly the energy of the scattered photons is similar to the incident photons and secondly, the resulting modes are only low-frequency waves due to velocities much smaller than the speed of light. Under these conditions, the derivation in § 2.2.3 is no longer valid since the thermal motion of the particles cannot be ignored. These conditions defined the possible plasma waves that can contribute to the shape of the scattered light. In this case, the two main low-frequency plasma waves are the electron plasma wave (EPW) due to free-electrons and ion acoustic waves (IAW) due to electrons closely bound to ions (the ‘Debye cloud’). For the low-frequency mode mediated by the collective motion of the electrons, the dispersion relation is given by the Bohm-Gross relation:

$$\omega^2 = \omega_{pe}^2 + \frac{3k_B T_e k^2}{m_e} \quad (2.22)$$

while the dispersion relation for the ion-acoustic waves follows the simple relation:

$$\omega = C_s k \quad (2.23)$$

The latter is derived by considering only sound waves with velocity $C_s = \sqrt{P/\rho}$ where P is the thermal pressure and ρ is the mass density. In weakly coupled plasmas such as those produced at MAGPIE, the sound speed can be expressed as $C_s = \sqrt{(ZT_e + T_i)/m_i}$ assuming quasi-neutrality ($n_e = Zn_i$) and ideal gas ($P = n_e k_B T_e + n_i k_B T_i$). This mode represents the collective behaviour of the electrons constrained by the inertia of the ions acting within the characteristic Debye length:

$$\lambda_D = \sqrt{\frac{k_B T}{4\pi n_e e^2}} \quad (2.24)$$

In addition to these two low-frequency modes, a third contribution is represented by the thermal motion of the electrons, resulting in non-collective fluctuations. Among these, the dominant mode depends solely on whether the wavelength of the scattering wave, $\lambda = 2\pi/k$, is longer or shorter than the characteristic length λ_D . For wavelengths much shorter than Debye length, the electrons appear free resulting in a scattering spectrum dominated by electron thermal motion. If instead, the scattering wavelength is longer than λ_D the bulk motion of the electrons is correlated with the

ion motion so both electron plasma wave and ion acoustic wave dominate the scattered spectrum. These two distinct regimes are referred to as non-collective and collective Thomson scattering which are characterised by the α - parameter, $\alpha = 1/k\lambda_D$. $\alpha \ll 1$ represents the non-collective regime while $\alpha \geq 1$ corresponds to the collective Thomson scattering regime. Under the typical MAGPIE plasma regimes, the characteristic Debye length is $\lambda_D \sim 0.2$ nm, which leads to $\alpha \simeq 3 \gg 1$ for a laser wavelength $\lambda = 532$ nm.

The spectrum of the scattered light due to Thomson scattering in the collective regime depends on density fluctuations which contain information about properties of the plasma such as electron and ion temperatures, the effective charge state of the plasma, the magnetic field, and electron density. By knowing the theoretical Thomson scattering spectral density function, it is possible to infer directly from the scattered light the plasma state. The derivation of the spectral density function, also known as the form factor, is complex and not presented here, limiting the discussion to its assumptions and characteristics. For more details refer to [84]. The spectral density function used in this study is the following:

$$S(k, \omega) = \frac{2\pi}{k} \left| 1 - \frac{\chi_e}{\epsilon} \right|^2 f_{e,0} \left(\frac{\omega}{k} \right) + \frac{2\pi Z}{k} \left| \frac{\chi_e}{\epsilon} \right|^2 f_{i,0} \left(\frac{\omega}{k} \right) \quad (2.25)$$

$$\epsilon = 1 + \chi_e + \chi_i \quad (2.26)$$

where, $f_{e,0}$, $f_{i,0}$ are the distribution functions of electrons and ions, respectively with $\chi_{e/i}$ the respective susceptibility. This form factor represents the Thomson scattering spectrum under the hypothesis of unmagnetised and collisionless plasmas. This may seem unintuitive given the plasma regimes produced at MAGPIE, however, it can be demonstrated that both assumptions are valid for the plasmas studied here. Specifically, the unmagnetised assumption is supported by the fact that its contribution to the electrons is small for most of the experimental geometry presented in this thesis, as discussed in section §2.2.3, while ions are typically unmagnetised. Similarly, any collisional effects happen at spatial scales much larger than the scattering interaction, $k\lambda_{ei} \simeq 10 \gg 1$ (with $\lambda_{ei} \simeq 1 \mu m$ electron-ion collisional mean free path), thus the plasma can be considered collisionless for Thomson Scattering [78]. The theoretical spectrum (2.25) consists of two terms: a term dependent on both the

non-collective thermal motion of the electrons and the collective electron plasma wave ('electron-component'); a term proportional to the ion distribution function describing the collective scattering of the electrons correlated to the ion motion ('ion-component'). In all the experiments presented in this work, Thomson scattering diagnostic operates in the collective regime with $\alpha > 1$. This means that the spectrum is only dominated by the ion component while the electron component is negligible.

Although in principle the ion distribution $f_{i,0}$ can assume any distribution function, in practice it is difficult to experimentally distinguish different distributions since there is no unique relation between plasma parameters and distribution functions. An appropriate assumption, in this context, is represented by a Maxwellian distribution expressed in terms of the velocity v :

$$f_{i,0}(v) = \sqrt{\frac{1}{2\pi v_{th,e|i}^2}} e^{\frac{-v^2}{v_{th,e|i}^2}} \quad (2.27)$$

where the thermal velocity is defined as $v_{th,e|i} = \sqrt{2k_B T_{i|e}/m_{i|e}}$. The plasma susceptibilities are given by the following expressions:

$$\chi_i = -\frac{\alpha^2}{2} \frac{\bar{Z} T_e}{T_i} Z' \left(\frac{\omega}{k v_{th,i}} \right) \quad (2.28)$$

$$\chi_e = -\frac{\alpha^2}{2} Z' (x_e) \quad (2.29)$$

$$x_{i|e} = \frac{\omega}{k v_{th,i|e}} \quad (2.30)$$

$Z'(x_{i|e})$ is the derivative of the plasma dispersion relations, (2.22) and (2.23), with respect to $x_{i|e}$. It is possible to express the above equations so that the Doppler shift due to the bulk motion of the plasma $V_{i|e}$ is included:

$$x_i = \frac{\omega/k + V_i}{k v_{th,i}} \quad (2.31)$$

$$x_e = \frac{\omega/k + V_i + V_e}{k v_{th,e}} \quad (2.32)$$

The form of x_e is chosen so that when the electrons move together with the ions, $V_e = 0$, where V_e is the relative velocity of the electrons with respect to the ions. When this relative velocity is not zero, the ion component is asymmetric due to Landau damping. This effect is stronger depending on the speed of the waves relative to the electron thermal velocity.

Thomson scattering diagnostic at MAGPIE

The Thomson scattering setup used at MAGPIE, as shown in Fig. 2.9, involves an Nd:YAG seed laser operating at 1064 nm with an 8 ns pulse length. The second harmonic (532 nm) obtained from the passage of the seed beam through a KPD crystal is then amplified to 3 J. The beam is focused by a $f = 1000$ mm lens to a focal spot at the centre of the experimental vacuum chamber, with a beam waist of 200 μm . A typical Thomson scattering setup involves the use of two collection lenses to image the scattered light from the chamber onto two linear arrays of optical fibre bundles (for simplicity, only one fibre bundle is shown in Fig. 2.9). The positions of the collection lenses and fibre bundles are determined based on the wanted measurements, defining the scattering vector $\mathbf{k}$. The collection optics in this setup are positioned at an angle θ relative to the incoming laser beam, and the scattering volumes for each fibre bundle are overlapped, allowing for two measurements of the scattered spectrum from different scattering vectors per scattering volume. The projection of the fibres onto the laser path through the imaging lens defines the location of the corresponding scattering volume, which represents the volume of plasma from which the scattered light is collected. Various fibre bundles are available at MAGPIE from 7 up to 44 fibres with a core diameter of 100 μm or 200 μm . The fibres are evenly spaced over a spatial extent between a minimum of 1.5 mm up to a maximum of 10.75 depending on the considered fibre bundle, the full inventory and fibre bundle specifications are reported in [61]. Based on the design of the imaging system, the spatial range of the measurements can be longer than the nominal fibre spacing, as this is determined by the magnification of the imaging system M , the beam waist along $\mathbf{k}_{\text{out}}$ and the projected size of the fibre core (185 μm diameter). The light collected by each fibre is coupled to a spectrometer (Andor SR 5001-A) equipped with an intensified CCD. The spectrometer CCD is gated with a 4 ns exposure, setting the temporal resolution of the experiment.

Various measures are taken to ensure a high signal-to-noise ratio, including baffling inside and outside the chamber to prevent laser light scattered by hardware or chamber windows from reaching the fibre bundles. Additionally, a narrow time gate on the spectrometer CCD minimises the impact of self-emission by the plasma. Temporal offsetting is employed for the 532 nm interferometry laser to prevent any of its light from reaching the spectrometer. Despite these measures, the measured spectra usually contain a notable contribution from background self-emission and laser light scattered by

experimental hardware. For alignment purposes, a metallic object in the form of a 100 μm diameter

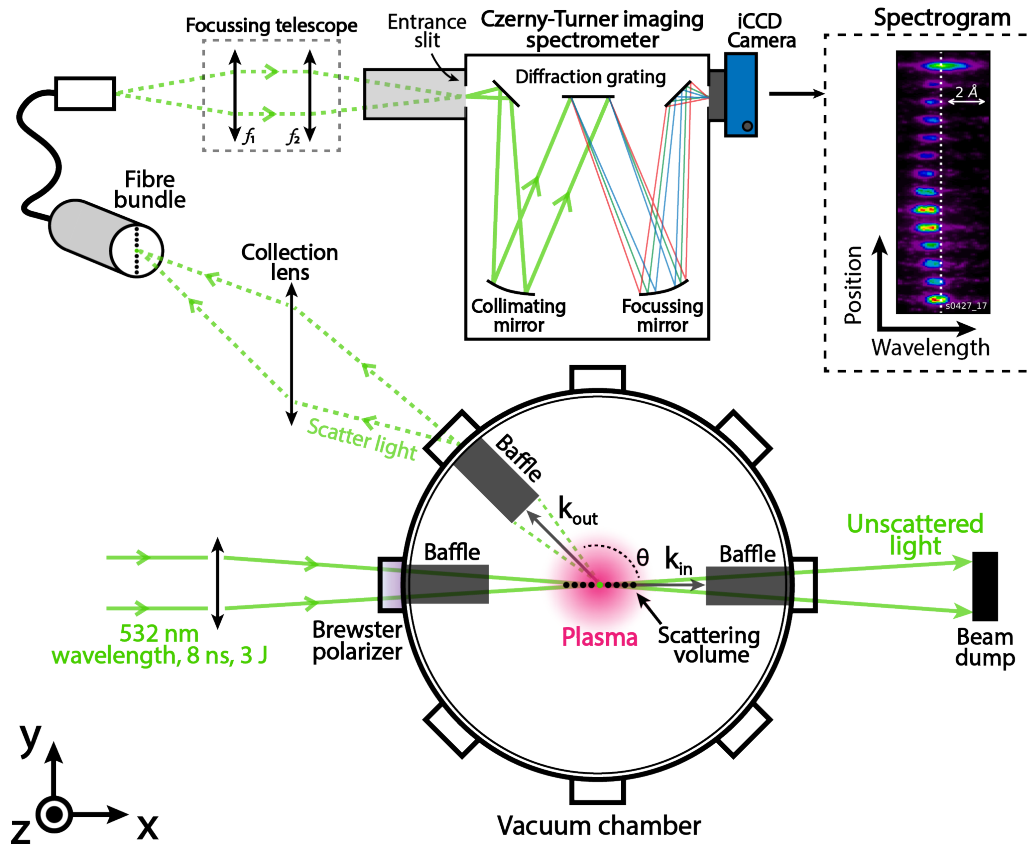


Figure 2.9: Schematic of the optical Thomson scattering setup. The probing laser (solid green) is focused on a spot inside the plasma by a lens. A single fibre bundle is shown in the illustration. Collection lenses are placed outside the vacuum chamber, imaging the scattered light (dashed line) onto the fibre bundles.

acupuncture needle is mounted on a three-axis vacuum-compatible micrometre stage as a proxy of a scattering volume. Placing the end of the needle within the laser beam produces a ring of scattered light which can be used for alignment. By adjusting the position of the needle within the laser beam in order to maximise the light that is incident on each fibre, the position along the beam is determined from which each fibre will record the Thomson scattering spectrum. This alignment technique offers an accuracy of approximately 100 μm . The position of each scattering volume in the measurement plane is accurately determined by pre-shot images of the light scattered by the needle from both side-on and end-on laser interferometry. These fibre positions can be then used to determine the electron density for each scattered spectrum by overlaying them onto the interferogram. This information is crucial for accurately fitting the Thomson scattering data.

Spectrum analysis of the ion acoustic feature

As already discussed previously, the optical Thomson scattering technique implemented at MAG-PIE operates in a collective regime with $\alpha > 1$, resolving the ion acoustic feature of the scattered light spectrum. In this regime, the form factor of the scattered spectrum is strongly sensitive to the plasma parameters, specifically to n_e , T_i , ZT_e , v_i and v_e . In addition, self-emission from the plasma and laser light reflections from the hardware can contribute ultimately to changes in the signals recorded by the spectrometer. Fitting the raw spectra to infer the state of the plasma is therefore non-trivial and an in-house code has been developed to assist in this analysis. A complete description of the code can be found in ([73], p. 67) while here just a general overview is provided. The ultimate aim of the analysis is to find the right combination of free parameters of the theoretical form factor eq. (2.25) such that it fits accordingly to the raw spectrum collected during the experiment. It is important to note that this process is not automatic and the free plasma parameters are manually chosen to find the best fit.

To help reduce uncertainties during the analysis, two fibre bundles are usually employed with overlapping scattering volumes so that the scattered light for a given spatial location is recorded twice, constraining the choice of T_i and ZT_e so that both spectra have the same values. Moreover, electron density n_e is usually known from the interferometry measurements, reducing, even more, the possible degrees of freedom, hence the source of errors. An example of spectral density function together with the convolved scattering signal recorded by the spectrometer is illustrated in Fig. 2.10(a).

The example illustrates in dashed black line the form factor $S(k, \omega)$ for a silicon-based plasma ($A = 28$) with $T_i = 25$ eV, $T_e = 25$ eV, $n_e = 5 \times 10^{18} \text{ cm}^{-3}$, $Z = 5$, $v_i = v_e = 0$ km/s. In this case, an angle $\theta = 90^\circ$ has been assumed between the scattering vector and the laser beam. The spectrum shape consists of two peaks, representing the ion acoustic waves travelling parallel and anti-parallel to the $\mathbf{k}$ vector. The recorded scattered spectrum by the spectrometer is shown in the solid cyan line as the result of the convolution between the theoretical form factor and the response of the spectrometer (solid grey line). Although peaks are still resolvable from the convoluted signal, the signal shape is significantly affected, leading to uncertainties in the measurements. The response function of the spectrometer is well approximated by a Gaussian distribution with a non-negligible FWHM of $\simeq 0.3 \text{ \AA}$. This finite broadening can limit the measurement capability of the diagnostic to resolve the double-peak feature in

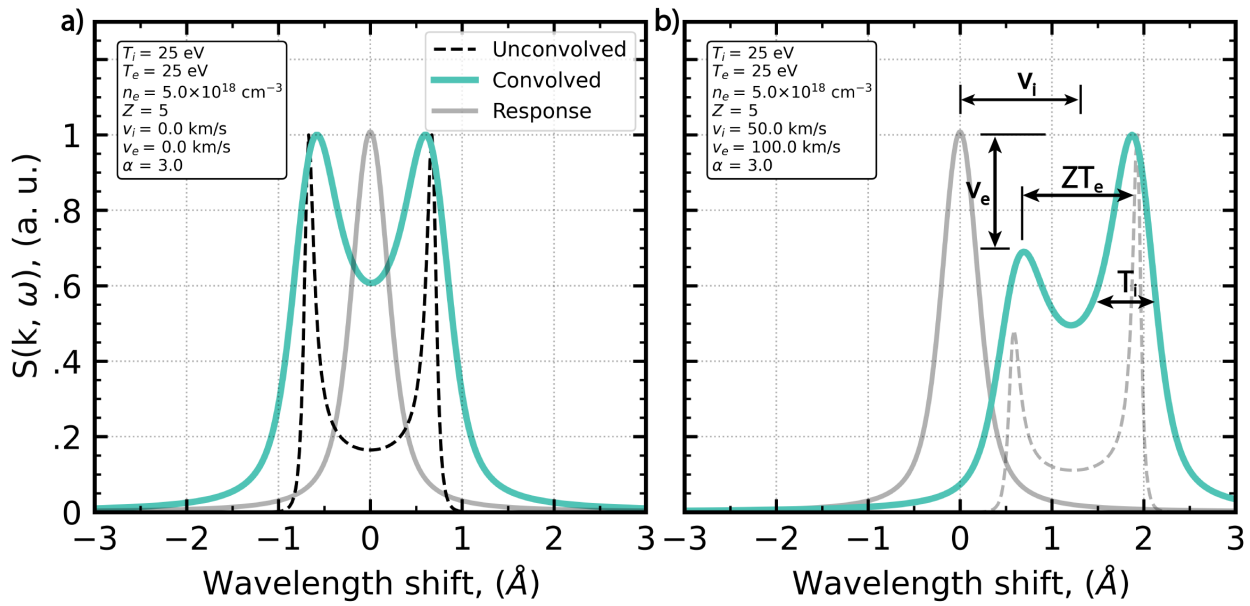


Figure 2.10: (a) Example of Thomson scattering spectrum of the ion-acoustic feature. The recorded Thomson scattering spectrum is the result of the convolution between the Thomson scattering form factor and the response function of the spectrometer. (b) The measurable plasma parameters from the Thomson scattering spectrum of the ion acoustic feature.

dense and cold plasmas, particularly when applied to high- Z elements such as tungsten (see Chapter 4). Fig. 2.10(b) shows a similar example, highlighting also the contribution of the electron drift velocity relative to the ions, v_e , and the ion bulk velocity v_i to the overall spectrum. From the plot, the thermal properties of the plasma act on the broadening and separation of the peaks, which are associated with T_i and the product ZT_e , respectively, while the bulk velocity v_i induces a Doppler shift of the whole spectral signal. It is important to note that the ion component of the form factor $S(k, \omega)$ is sensitive to the product ZT_e and not to the two parameters, separately. This means that, an estimate of the electron temperature T_e can only be achieved by assuming an ionisation model, $Z(T_e, n_e)$. In this case, look-up tables from two atomic codes are typically used, the non-local thermodynamic equilibrium code SPK [85] and FLYCHK [52]. The asymmetry of the peaks is affected by the relative velocity v_e , which in some cases can be used to estimate the plasma current. From these considerations, it can be stated that Thomson scattering in a collective regime allows to infer a broad range of plasma parameters, especially when combined with interferometry data.

2.3 Platforms for HED plasma experiments

In this thesis, two platforms have been used to explore reverse shocks in high-energy-density plasma experiments. The first one is based upon already established research in MHD shocks, consisting of an exploding wire array where plasma flows are accelerated radially outwards due to the Lorentz force. The second platform, instead, is a novel design built to overcome some of the limitations of the traditional exploding wire arrays. The idea here is to use an imploding wire array as an X-ray source to radiatively ablate solid targets. In this section, both platforms are presented, giving a general overview of the physical mechanism, while leaving details about the geometry later in the experimental Chapters.

2.3.1 The inverse Z-pinch wire array platform

The first experimental campaign employed an inverse (or exploding) wire array z-pinch as a plasma source, consisting of metal wires cylindrically arranged around a central cathode. Inverse wire arrays have been used at MAGPIE for more than a decade, studying various research topics around the laboratory of astrophysics and HED plasmas [59, 60, 86–90]. For the sake of brevity, only a brief discussion is intended while for more detailed information the reader can refer to [91].

A CAD drawing of the platform is shown in Fig. 2.11. The hardware consists of an anode structure formed by an anode ring supported by an anode plate sustained by 8 return posts attached to the vacuum chamber breadboard through a stainless steel disk. The central cathode structure is connected directly to the central MITL by a hollow support which sustains the cathode disk of the array. During the experiment, the current pulse from the generator flows from the anode through the thin metal wires to the cathode disk, before descending through the cathode post. A magnetic field fills the gap between the wires and the cathode post, generating a strong magnetic pressure within the wire array. Ohmic heating of the wires causes a plasma sheath to form around the cooler wire cores which is then accelerated radially outward under the presence of the $\mathbf{J} \times \mathbf{B}$. As the plasma moves away from the wires, new plasma forms on the wire surfaces, resulting in a continuous ejection of plasma from each wire. When the wires are closely spaced, the outflows from adjacent wires merge, leading to a smooth

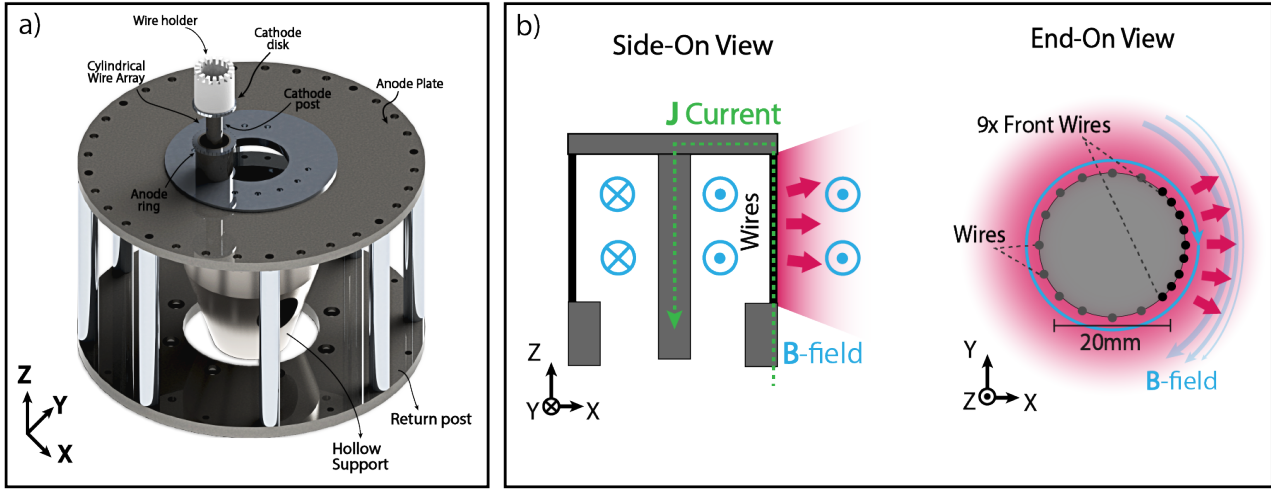


Figure 2.11: (a) CAD drawing of an exploding wire array platform used at MAGPIE. The hollow support is mounted on the central cathode of MITL, while the bottom anode plate is attached to the breadboard of the vacuum chamber. (b) Schematic representation of both side-on and end-on views of the inverse wire array configuration.

plasma outflow.

The main advantage offered by this configuration over standard imploding wire arrays is the absence of any obstruction in the plasma field of view. As the plasma flow is radially outward, the diagnostic lines of sight are not interfered with by the wire cores or surrounding hardware. A good mathematical description of the plasma outflow is given by the same rocket model initially developed for describing wire ablation in an imploding wire array z-pinch [59]. This model describes the rate of mass ablation from the wires by balancing the $\mathbf{J} \times \mathbf{B}$ force with the radial momentum [92]:

$$v_{abl} \frac{dm}{dt} = -\frac{\mu_0 J^2}{4\pi R} \quad (2.33)$$

Where v_{abl} is the ablation velocity and R is the radius of the wire array. Fig. 2.12 illustrates the magnetic field profile surrounding the wire array in its early stages [91]. This magnetic field is primarily azimuthal between the cathode and the wires, with minimal field outside the array. Each current-carrying wire generates a local magnetic field superimposed on the global magnetic field surrounding the cathode, leading to small regions of private flux around each wire. Initially, the ablated plasma remains close to the wires, resulting in a magnetic Reynolds number less than 1, $Re_M = vL/\eta < 1$, due to the small characteristic length, L . However, as the plasma moves away from

the wire surface and L increases, leading to the growth of Re_M . Consequently, some of the magnetic field initially inside the array becomes frozen into the plasma, creating an expanding plasma outflow with a naturally embedded magnetic field.

Prior experiments have shown that the plasma outflow produced from various wire materials, such as aluminium and copper, is supersonic and super-Alfvénic. The advected magnetic field is dynamically significant, with magnetic pressure comparable to the thermal pressure but smaller than the ram pressure. The outflow velocity reaches approximately 100 km/s, and the advected magnetic field strength ranges from 1 to 5 T. Typical electron densities are in the range of 10^{17} to 10^{18} cm $^{-3}$. These characteristics make wire arrays ideal for studying shocks in a magnetized plasma, as a super-fast-magnetosonic flow will form a shock around a stationary object placed within the flow.

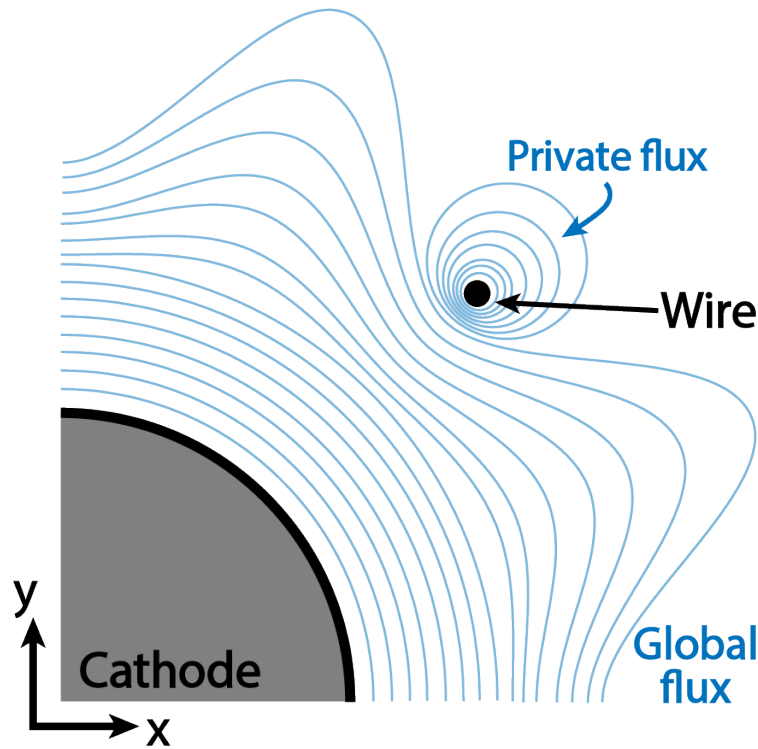


Figure 2.12: Magnetic flux contours in an exploding wire array just after the current pulse starts. A global flux is formed around the central cathode while superimposed private fields are generated around wire cores from each single current stream.

For experiments requiring a sustained plasma flow, such as those presented here, the wire arrays are usually over-massed. This involves using thicker wires than conventional wire arrays or a sufficient number of wires so that any explosions occur during the experimental timeframe ($\simeq 500$ ns). Unlike a standard imploding wire array, where the initial ablation phase lasts until about 50% of the total

ablated wire mass, the over-massed wire array (in an inverse configuration) sustains a steady stream of plasma throughout the entire drive current duration. Although these wire arrays still eventually explode later in the experiment, this only occurs after the current drive has ended which is distinct from the one observed in standard wire arrays.

2.3.2 The X-Ray driven platform

In contrast to the previously described inverse wire array, the second experimental platform used in this study involves X-ray ablation of solid targets using a standard Z-pinch wire-array imploder as represented in Fig. 2.13(a). Here, the plasma source does not originate directly from the wire array outflow; instead, it arises from the ablation of small solid targets (e.g. silicon, carbon, etc.). A detailed schematic of the experimental platform is shown in Fig. 2.13. The cathode structure, comprising the cathode plate and cathode disk, is mounted on the central conductor of the MITL (Magnetic Insulated Transmission Line) using a hollow support. To enable end-on probing of targets along the z-axis, the z-pinch array axis is offset from the chamber centre by 30 mm. The hollow support's cathode plate, top, and side have holes that facilitate directing the end-on laser probe into the imaging optics through a 45° mirror. At a height of 40 mm, an anode disk is attached to the return structure, and the 16 aluminium wires are hung from it, making contact with the cathode disk below. Grooves on both disks secure the wires in place. The anode structure consists of four return stand-off posts connected to the anode plate, which is then grounded by eight return posts to the chamber breadboard. The imploding wire array comprises 16 evenly spaced aluminium wires arranged cylindrically, with a diameter of 16 mm and a height of 40 mm. Each wire has a thickness of 20 μm and connects the central cathode to the outer anode structure surrounding the wire array. The initial selection of array parameters aims to maximize the total X-ray yield while minimizing experimental complexity. For the experiments presented in this thesis, a setup similar to the one described by [67] is employed, with some key differences. Additional plasma flow sources are introduced to investigate the interaction of two counter-streaming colliding flows. The orientation of the ambient magnetic field relative to the plasma flows is also explored. The number of wires is reduced from 32 to 16 while increasing the mass nearly twice to ensure a longer X-ray pulse duration while avoiding sudden wire breakdown. Exploring different ablating target

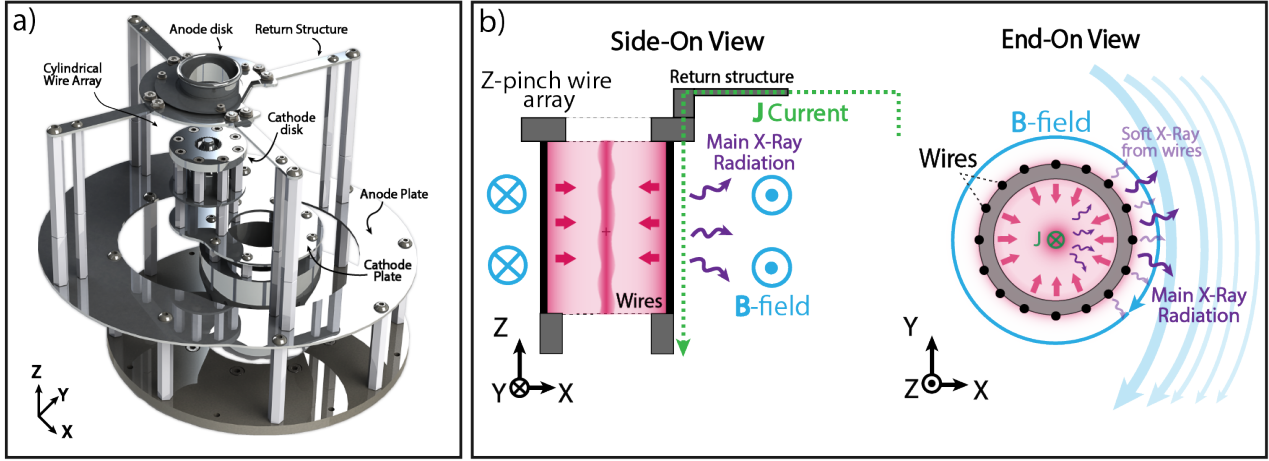


Figure 2.13: (a) 3D render of the X-ray driven platform. X-ray radiation is produced from the implosion of a cylindrical wire array off-axis relative to the central cathode. (b) Schematic representation of both side-on and end-on views of the imploding wire array configuration.

materials could potentially alter the spectral characteristics of the X-ray pulse and the overall radiation yield, providing advantages for future experiments.

As the current pulse begins, the wires undergo Ohmic heating forming the wire core, presented in the previous section. Despite being relatively cold and dense, the wire core still provides a path for the current from the generator. A coronal plasma forms around these wire cores and a fraction of the generator current flows into this plasma, initiating the ablation phase. During the experiment, the magnetic field configuration is similar to the one shown in Fig. 2.12. The advection of the magnetic field leads to a switch to a global magnetic field topology with no private flux, where roughly azimuthal field lines cross through each wire core, resulting in a $\mathbf{J} \times \mathbf{B}$ force radially inward on the coronal plasma [93]. This array implosion dynamic results in the production of X-ray radiation as a two-step process, an early X-ray precursor produced from the early ablation phase of the wires followed by a main X-ray pulse from the full array implosion [92, 94]. As already mentioned, the amount of X-ray energy produced and the X-ray pulse shape is determined by the combination of wire material, number of wires and geometrical size of the array, as described in [92, 93].

In the experiments reported in this thesis, the contribution of the X-ray precursor is negligible to the ablation of the solid targets with the plasma flows driven only by the radiation produced from the "Snowplow" implosion phase. An example of a wire array implosion during one of the X-ray driven experiments is reported in Fig. 2.14. The Z-pinch wire array was fielded with an XUV fast-frame

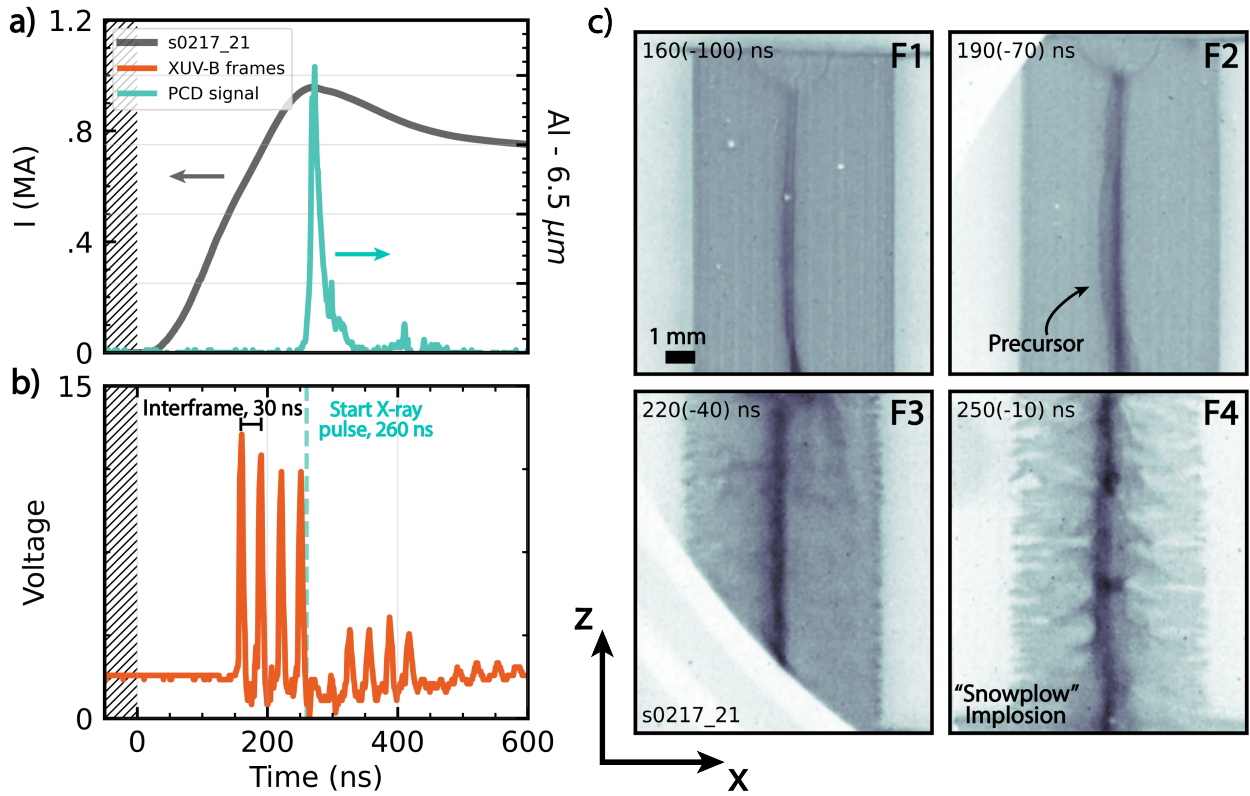


Figure 2.14: (a) The current profile and the signal of an Al-filtered PCD detector show the start of the main X-ray pulse. (b) Timing signal of the 4-frame XUV fast-frame camera. The fourth peak at 260 ns after the current pulse corresponds to the "snowplow" implosion phase (frame F4), which produces the main X-ray pulse. (c) Side-on XUV fast-frame pin-hole imaging (photon energy > 40 eV, 5 ns exposure) of the wire array implosion phase during the s0217_21 experiment.

pin-hole camera from the side-on view together with an Al-filtered PCD detector to record the produced X-ray pulse. The main X-ray pulse is produced at the current peak of the generator, as shown in Fig. 2.14(a). This corresponds to the stagnation of the wire array implosion as represented by the last frame (F4) in the XUV pin-hole camera sequence of Fig. 2.14(c). As will be discussed further in Chapter 5, the moment of the "Snowplow" implosion represents the start of plasma flows from the solid targets, thus the reference time in the X-ray driven experiments.

Chapter 3

Imaging Refractometry: A novel Refractive-Based Diagnostic

Refractometry-based diagnostics, such as Shadowgraphy and schlieren imaging, have been demonstrated to be a suitable tool for studying qualitative shocks and density fluctuations in High Energy Density plasmas. In both cases, laser probing rays are deflected by plasma density structures, acting as natural lenses which results in intensity variations at the image plane. Although often used as quantitative methods, the loss of a one-to-one relationship between the image on the detector and the object due to "caustic" effects makes these techniques limited to a narrow range of gradients and spatial scales [79].

To complement these conventional techniques, a new hybrid diagnostic called the Imaging Refractometer [62] has been developed at MAGPIE with the aim of overcoming some of the limitations of traditional refractive-based diagnostics when properties of density fluctuations are homogeneous in one direction. The imaging refractometer uses a combination of cylindrical and spherical lenses to image the distribution of deflection angles of a scattered laser beam while it propagates through a plasma medium. This results in a hybrid image where on the vertical axis are deflection angles and on the horizontal axis are spatial positions, reaching over 500 in the dynamic range, which is far beyond conventional methods. In addition, integrating over the spatial axis results in an increased signal-to-noise ratio without the formation of caustic effects.

In this chapter, the novel imaging refractometer technique is first introduced, describing both physical principle and optical arrangement. A synthetic tool based on geometrical optics approximation has been developed to further understand the output of this diagnostic and it is then discussed. Finally, experimental data collected from the imaging refractometer applied to a simple test case is presented, allowing for validation of the diagnostic.

3.1 Optical configuration

The imaging refractometer shares the same physical principles of conventional refractive-based diagnostic, such as shadowgraphy or schlieren, presented in the previous chapter, 2. In an inhomogeneous medium, the laser beam light is deflected by the presence of density gradients normal to the propagation direction according to eq. 2.16. This distribution of deflected angles is directly related to the line-integrated transverse density gradients, dependent on the density structures in the plasma. The novelty behind the imaging refractometer technique is to measure the distribution of deflection angles, so that information about density gradients and structures within the plasma electron density can be obtained.

The imaging system shown in Fig. 3.1(a) uses physical optics to simultaneously capture a one-dimensional spatially resolved deflection angle distribution and a conventional shadowgraphy image of the probing beam. The system employs the long-pulse arm of the Nd:glass CERBERUS laser ($\lambda_0 = 1053$ nm, 100 mJ, 1 ns) to ensure high reproducibility of the intensity profile and low beam divergence, < 0.05 mrad. This ensures that any intensity variations are due to deflecting light rather than changes in the light source, while small angular divergence ensures that the light beam can be well represented as a single Fourier component at the sensor plane, resulting in a very narrow line. Moreover, the short coherence length (< 1 mm) and the infrared characteristic wavelength reduce interference effects while maximising the system sensitivity. In the experiments presented in this thesis, the laser beam passes through the plasma from side-on, along the y-axis, imaging the xz-plane.

The optical configuration consists of a one-lens and a two-lens imaging system, with the probe light directed to the shadowgraphy and refractometer lines using a 50:50 beam splitter. The shadowgraphy

line represented in Fig. 3.1 (b) is a classical shadowgraphy system composed by a single spherical lens $f_1 = 200$ mm (50 mm diameter Thorlabs Achromatic doublet AC508-200-C-ML) shared among the two arms, forming a 2D shadowgraphy image of the plasma object at the camera plane, C1 (Canon 500D). For the imaging refractometer, a second optic is placed at a distance $3f_1$. This consists of a hybrid lens made of a cylindrical lens and a spherical lens placed back-to-back with a total focal length, $f_3 = 400/3$ mm, which allows imaging of the Fourier plane along the vertical axis of camera C2, while maintaining spatial positions along the horizontal axis. An ATIK 383L+ camera was used with 3448×2574 pixels, 18×13.5 mm, monochrome Kodak KAF-8300, 16-bit depth CCD. The high bit-depth and overall low noise ensure the highest achievable dynamic range of the system when imaging deflected rays. The exposure time of both cameras is kept longer than the duration of each experiment so that the temporal resolution is set by the laser pulse duration, 1 ns. An example of a

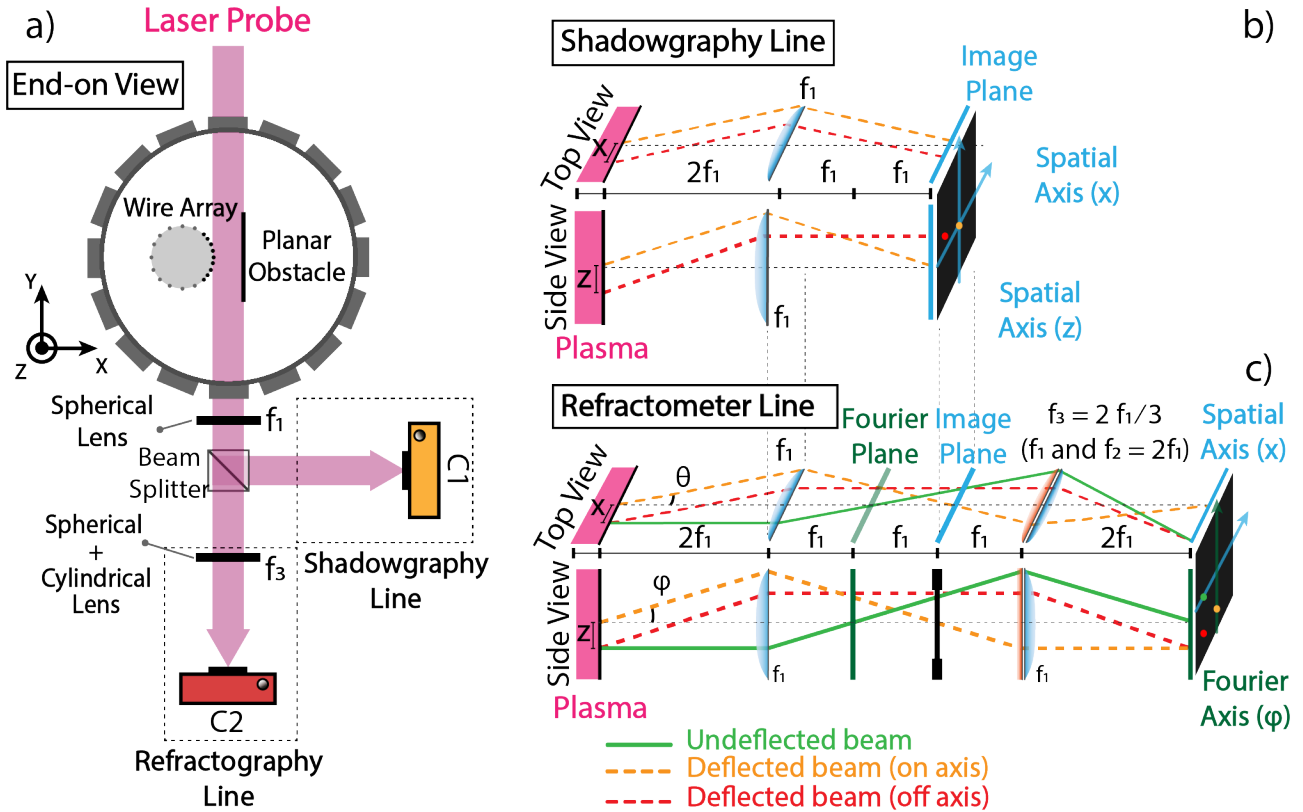


Figure 3.1: End-on view shadowgraphy and imaging refractometer system. Ray diagrams for both shadowgraphy and refractometer.

ray diagram for the imaging refractometer is shown in Fig. 3.1(c). After the first spherical lens, f_1 , two planes are formed, an Image plane and a Fourier plane, containing respectively the spatial and angular distribution of the rays after passing through the plasma object. The second optic consists

of a spherical lens, f_1 , and a cylindrical lens, $f_2 = 2f_1$, which focuses light differently along two orthogonal axes. The green line represents a beam of collimated light that is propagating parallel to the optical axis. This ray intersects the centre of the Fourier plane, and it crosses the image plane at the same distance from the optical axis as it originated. After passing through the second lens, this ray arrives at the optical axis in the Fourier direction (at an angle of $\phi = 0$), and on the imaging axis, it reaches the detector at twice its initial distance, which corresponds to a magnification of $M = 2$. The orange and red rays represent a collimated laser beam that has been deflected by an angle relative to the optical axis. These rays intersect the first Fourier plane off-axis and reach the image plane with a magnification of $M = 1$. The second lens images both rays to the same off-axis location in the Fourier axis, which indicates that the ray exited the plasma with an angle of $\phi \neq 0$. Both rays are imaged with a magnification of $M = 2$ onto the imaging axis. The resulting image recorded by camera C2 will thus represent the distribution of deflection angles $\phi(x)$, along the axis ϕ , for various spatial positions x . Examples of output from this imaging system are shown later in the section.

3.2 Synthetic Diagnostics Tool

Along with the development of the imaging refractometer diagnostic, a synthetic tool based on a custom ray tracer has been developed to demonstrate the capability of this new diagnostic to quantify plasma properties, such as electron density gradients. This numerical tool allows to generate synthetic refractograms based on a provided three-dimensional plasma volume. This has two main advantages: study the response of the diagnostic under simple test cases and assess uncertainties and errors in the data analysis, showing limitations and possible improvements of the diagnostic itself.

The synthetic diagnostic code is made of two parts: a laser ray tracing algorithm that initialises and propagates rays through the plasma volume and an optical system module which instead propagates the deflected rays through the optics up to the image plane at the detector. It is important to note that in both cases, geometrical optics approximation is assumed, considering only the cases of small deflection angles. Under this condition, displacements in x and z are negligible so that only gradients in the xz -plane are important. Moreover, rays propagate through the plasma domain following the

simple geometric optics equations, which can be evaluated at discrete positions across the plasma volume, as illustrated in Fig. 3.2(Left):

$$\frac{\Delta \vec{v}}{\Delta t} = -\nabla \left(\frac{c^2 n_e}{2 n_c} \right) \quad (3.1)$$

$$\frac{\Delta s}{\Delta t} = \vec{v} \quad (3.2)$$

First, a bundle of rays with uniform distribution is initialised, representing the probe beam incident to the plasma slab, X_0 . The beam is assumed to be perfectly parallel to the propagation direction (y-axis) with initial angles $\theta = 0^\circ$ and $\phi = 0^\circ$, relative to the x and z directions, respectively. Then, the electron density gradient across one propagation step along the laser path is computed by taking the central difference between each step size, as shown in Fig. 3.2(Right). Finally, using the discrete version of eq. 2.16, the deflection angles are updated along with the direction of the ray, becoming the initial state for the following iteration step. For more details refer to [95]. The result of the ray tracing algorithm

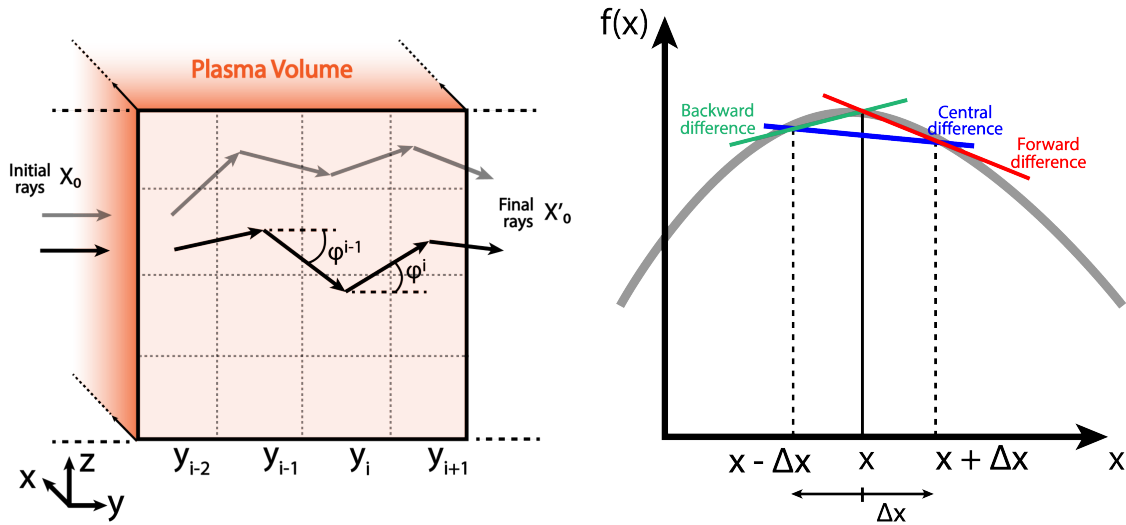


Figure 3.2: Illustration of ray propagation through a plasma slab. The ray propagates along the y-axis and it is deflected by the presence of density gradients on the transverse plane (xz-plane). The associate deflection angles are θ and ϕ , respectively for the components of density gradient, $\partial_x n_e(x, y_i, z)$, $\partial_z n_e(x, y_i, z)$. Illustration of the gradient approximation in a 1-dimensional function, $f(x)$. In this case, the gradient is estimated from the central difference between two discrete steps.

represents the final state of the laser beam at the exit of the plasma volume, X'_0 , before passing through the optical imaging system. Under paraxial approximation, the optical arrangement illustrated in Fig. 3.1(b) & (c) can be treated using standard ray-transfer matrix (RTM) techniques, [96]. At its core, the

RTM approach represents optical systems as a series of matrices, each corresponding to an optical element or interface. These matrices capture the transformation of ray parameters, such as positions and angles as light passes through the system. By cascading these matrices, the complete transfer function of the optical system can be obtained. In this case, all of the lenses are considered to be thin objects described by the following 4x4 matrices:

$$F_j(f_1, f_2) = \begin{bmatrix} 1 & 0 & 0 & 0 \\ -1/f_1 & 1 & 0 & 0 \\ 0 & 0 & 1 & 0 \\ 0 & 0 & -1/f_2 & 1 \end{bmatrix} \quad (3.3)$$

Where f_1 and f_2 are the two focal lengths respectively along the two directions, x and z. If we consider a single optical element, the associated matrix, $F_j(f_1, f_2)$, operates on the ray vectors representing the incoming ray light such that the calculated outgoing ray is given by $X'_{i+1} = F_j X'_i$. Extending the example across each optical element of the imaging system for both shadowgraphy and imaging refractometer lines, the final positions of the ray at the detector can be resolved as follows:

$$X'_0 = \begin{bmatrix} x \\ \theta \\ z \\ \phi \end{bmatrix} \rightarrow X_f = X'_0 \cdot F_{shadow|refract} \quad (3.4)$$

F corresponds to the transfer matrix of the optical system under consideration. The shadowgraphy line consists of a single spherical lens with equal focal lengths in both lens axes, $f_1 = f_2 = 200$ mm. The plasma object and the detector are at an equal distance $L = 2f_1 = 400$ mm from the lens. The ray propagation in a vacuum is represented by the 4x4 matrix $S(L)$ with equal diagonal elements:

$$S(L) = \begin{bmatrix} 1 & L & 0 & 0 \\ 0 & 1 & 0 & 0 \\ 0 & 0 & 1 & L \\ 0 & 0 & 0 & 1 \end{bmatrix} \quad (3.5)$$

This results in the following transfer matrices for the shadowgraphy and the imaging refractometer:

$$F_{shadow} = S(L) \cdot F_1(L/2, L/2) \cdot S(L) \quad (3.6)$$

$$F_{refract} = S(L) \cdot F_2(L/3, L/2) \cdot S(3L/2) \cdot F_1(L/2, L/2) \cdot S(L) \quad (3.7)$$

If we propagate the position of the ray X'_0 at the exit of the plasma through the imaging refractometer system, the position of the rays at the detector can be calculated simply as:

$$X'_0 = \begin{bmatrix} x \\ \theta \\ z \\ \phi \end{bmatrix} \rightarrow X_f = X'_0 \cdot F_{refract} = \begin{bmatrix} 2x \\ \theta/2 + 2x/f_1 \\ -f_1\phi \\ z/f_1 \end{bmatrix} \quad (3.8)$$

Under paraxial approximation, the image formed at the camera C2 of the imaging refractometer is therefore an hybrid image: on the horizontal axis of the CCD, rays are imaged with a magnification of 2 and they have no dependence on the initial angle, while on the vertical axis of the detector, the position of the rays are only dependent on the angle of the initial ray (ϕ) from the y-axis in the xz-plane. The location of the ray at the detector is $f_1\phi$, which defines the angular sensitivity of the detector. The resolution of the instrument is set by any imperfections in the optics, the initial divergence of the beam, the pixel size of the detector, and limitations due to diffraction.

3.3 Synthetic Test Case: Exponential Monochromatic Perturbation

To illustrate the power and utility of the imaging refractometer, a simple test case using the ray tracing code presented above was performed, generating synthetic shadowgraphy and refractometer images of a predefined electron density. In this case, a 10 mm cubic volume was initialised with the

following electron density perturbation on the xz-plane:

$$n_e(x, y, z) = n_{e,0} 10^{-x/s} [1 + \cos(2\pi z/L_z)] \quad (3.9)$$

Along the z-axis, the n_e perturbation is monochromatic with characteristic wavenumber, $k_0 = 2\pi/L_z$, while along x the electron density decays as a base 10 exponential law with rate constant, s. The electron density is assumed to be constant along the laser probe direction (y-axis). A visual representation of the imposed electron density perturbation is shown in Fig. 3.3. The periodic component of the electron density distribution is along the vertical spatial z, while the overall density exponentially decays along the horizontal spatial x. In the current configuration, the imaging refractometer is mostly sensitive to the z-component of the electron density gradient along the vertical axis. It is worth noting that shadowgraphy effects are still present in the x-direction, however here the density gradient is much smaller than in the z-direction, so such effects are negligible. A uniform distribution of 10^8 rays was

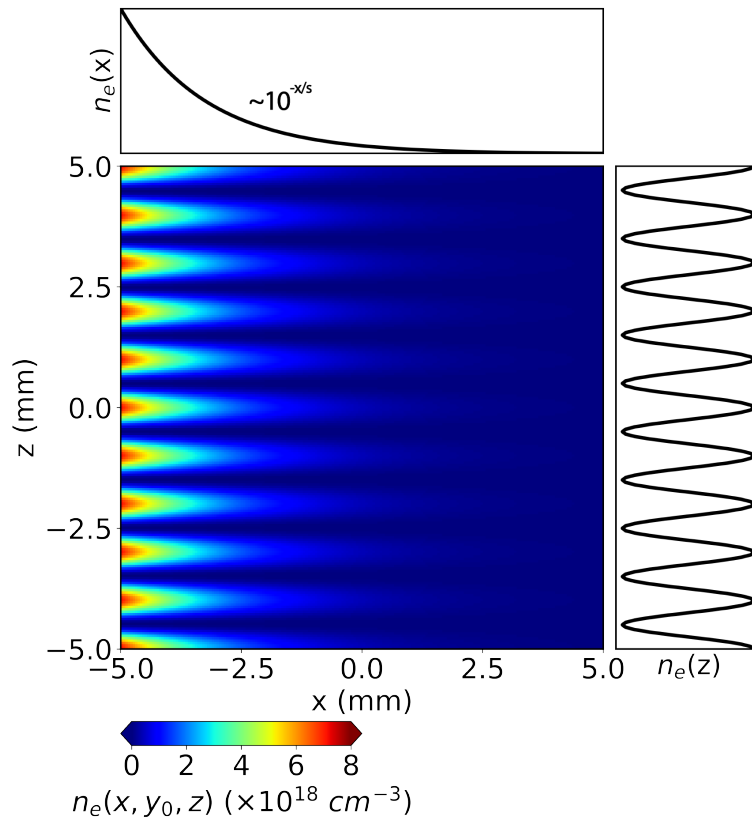


Figure 3.3: Example of a synthetic slice from the 3D electron density distribution eq. 3.9 with $s = 4$ mm, $L_z = 1$ mm and $n_{e,0} = 2 \times 10^{17} \text{ cm}^{-3}$ at fixed y_0 . The electron density profile along the x-direction decays exponentially, $n_e(x) \sim 10^{-x/s}$. In the z-direction, the electron density is periodic with $n_e(z) \sim \cos(2\pi z/L_z)$ while uniform along the y-direction.

propagated along the y-axis of the electron density cube. The ray tracing result was then converted into synthetic diagnostic images by a ray transfer matrix method assuming paraxial approximation (small deflection angles), taking into account the real optical arrangement used during the experiments.

Results are shown in Fig. 3.4(a) & (b), for the shadowgram and the refractogram respectively. Starting from the synthetic shadowgram, it is seen that the density distribution results in the formation of periodic caustics at around $x = -4$ mm where the density gradients are the largest. For increasing x , the overall intensity decreases until it reaches the diagnostic sensitivity. In this case, the laser probing light is deflected by the z-component of the electron density gradient which is then imaged at the detector plane. Imaging refractometer data represent hybrid images of optically Fourier-transformed

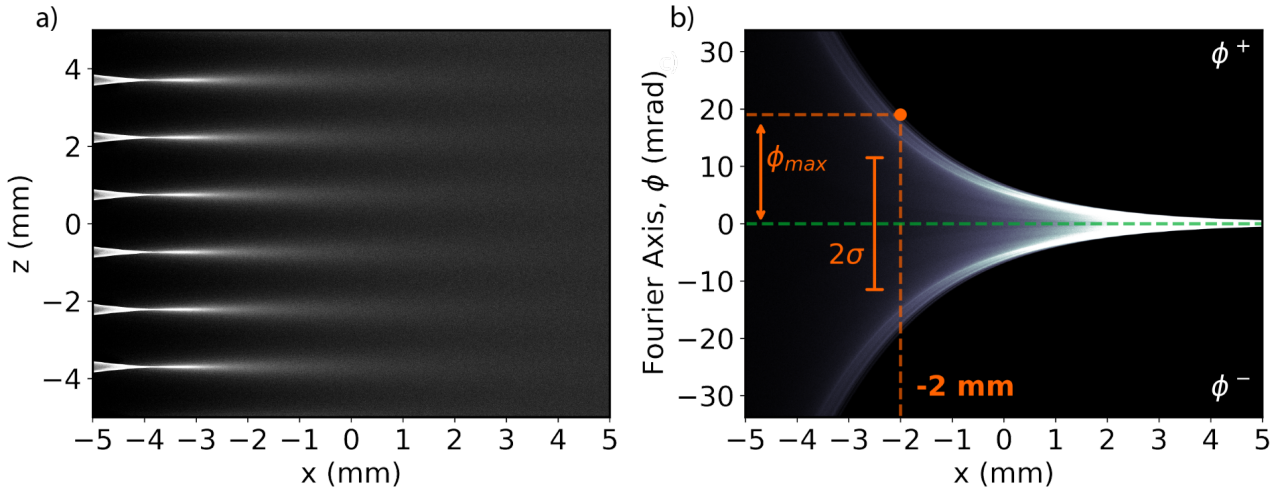


Figure 3.4: (a) and (b) are synthetic shadowgraphy and refractography images generated using a custom ray tracing code with 10^8 rays uniformly distributed on the xz -plane and propagating along the y -axis.

gradients along z on the vertical axis and spatial position along x on the horizontal axis. Since, in this case, electron density gradients are symmetrically distributed along the vertical axis, the equivalent refractogram results in two identical intensity profiles (ϕ^+ , ϕ^-) symmetric around $\phi = 0$ mrad axis. The maximum z -component of the density gradient can be estimated analytically by taking the first derivative of the electron density distribution, $n_e(x, y, z)$, calculated at $z = 0$ [78]:

$$\left(\frac{\partial n_e(x, y, z)}{\partial z} \right)_{\max} = \frac{2\pi}{L_z} 10^{-x/s} n_{e,0} \quad (3.10)$$

Knowing that the ray deflections are directly proportional to the gradient of the electron density

perturbation, the maximum deflection angle, ϕ_{max} , is:

$$\phi = \frac{1}{2} \int \frac{\nabla n_e}{n_{cr}} dy \Rightarrow \phi_{max} \approx \frac{\pi 10^{-x/s} n_{e,0} L_y}{n_{cr} L_z} \quad (3.11)$$

it is possible to further illustrate how the imaging refractometer operates by comparing the synthetic signal with its expected calculated value from eq. (3.11). For example, if we consider spatial position $x = -2$ mm, the maximum density gradient according to eq. (3.10) is $\simeq 4 \times 10^{19} \text{ cm}^{-3}/\text{cm}$ ($L_z = 1$ mm, $s = 4$ mm, $n_{e,0} = 2 \times 10^{17} \text{ cm}^{-3}$), which results in the maximum deflection angle of ϕ_{max} ($x = -2$ mm) $\simeq 20$ mrad. This value corresponds to the intensity profile on the refractogram (dashed orange line). If we now consider $x = 2$ mm, the resulting deflection angle will be 10 times lower, again in agreement with the positive intensity signal (ϕ^+). This is not surprising, since the assumed density gradients scale with a base 10 exponential along the x-axis, clearly shows that information about electron density gradients can be retrieved directly just from the signal ϕ^+ and symmetrically from ϕ^- .

In real experiments, however, the exact analytical expression of the electron density distribution is unknown and real refractometer data are often affected by light artefacts and noise that can compromise the image quality. In addition, the density gradient could vary along the probing beam direction, so a more general characterisation of the angular distribution can be considered by computing the standard deviation of spatially sampled (along the x-axis) vertical intensity distributions of the experimental refractograms. As observed directly from the raw image, stronger density gradients produce larger ray deflections, resulting in the broadening of the intensity distribution around its central value. Without any assumptions on the type of distribution, the standard deviation of the signal intensity represents a statistical description of angular deflections for a given spatial position, $\sigma_\phi(x)$. Higher-order moments, such as skewness and kurtosis, may also be used to analyse the distribution of angular deflections and possibly relate them to the spectrum of density fluctuations within a plasma. This thesis, however, will focus solely on the standard deviation and further statistical analysis will be explored in future works. The result of the analysis is shown in Fig. 3.5, averaged over $\Delta x = 0.2$ mm. Overall, the measured σ_ϕ follows a similar exponential law as the one predicted by the analytical expression of eq. (3.11) (blue line). From the exponential fit of σ_ϕ , the measured characteristic spatial decay (maximum deflection angle decay with $s = 4$ mm) is $s = 4.5 \pm 0.02$ mm, differing by $\approx 13\%$ from the actual value. In the

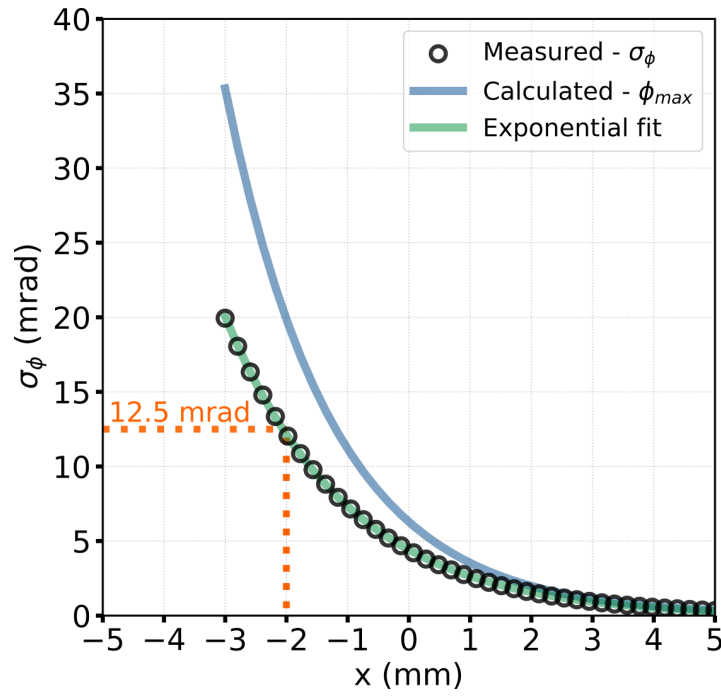


Figure 3.5: Comparison between the spatially resolved standard deviation, $\sigma_\phi(x)$, and the calculated maximum deflection angle, ϕ_{max} .

case of the synthetic test case, the standard deviation analysis results in a low estimate of the angular deflections due to the highly non-Gaussian distribution. Based on these results, however, the imaging refractometer successfully captures additional information about electron density gradients in plasma flows which are not resolvable from traditional shadowgraphy.

3.4 Experimental Validation - Shock Pattern formed by a Planar Grid

The synthetic study of the imaging refractometer applied to a simple electron density perturbation has demonstrated the potential of this new diagnostic to obtain quantitative insights on the plasma which were unavailable from the conventional shadowgraphy or schlieren techniques. To investigate and demonstrate the response of the imaging refractometer under laboratory conditions, an experiment employing an aluminium wire array and a high transparency grid of copper wires aligned with the magnetic field advected by the flow was performed, before applying this technique to a more complex system. As will be introduced in the next chapter, the same diagnostic arrangement was then applied

to study the structure of plasmas colliding with planar obstacles. The configuration of the wire arrays remained consistent throughout all experiments, as did the position of the arrays in relation to the obstacles – only the morphology of the target (wire grid replaced by planar obstacle) was changed in the final set of experiments.

In this section, the aluminium wire array experiment with a high-transparency grid is presented. A description of the experimental setup is provided together with optical self-emission images taken from the side-on view of the post-grid region. The shadowgraphy and imaging refractometer data are then provided along with an example of deflection angle analysis from real experiment data. The section concludes with a discussion of the results.

3.4.1 Experimental Overview

As already introduced in § 2.3.1, exploding wire arrays driven by a current pulse are a versatile system to generate highly diagnosable long-lasting plasma flows. In this experiment, the plasma material ablated from the wires is accelerated radially outwards by the presence of the $\mathbf{J} \times \mathbf{B}$ force, producing a quasi-uniform plasma stream which then collides against a highly transparent planar grid placed perpendicularly to the flow direction at 10 mm from the wire array, as illustrated in Fig. 3.6. The grid consisted of 16 evenly spaced copper rods with a diameter, $D = 50 \mu\text{m}$, and a spacing, $\delta_0 = 800 \mu\text{m}$, oriented parallel to the y-axis (which is the same orientation as the advected magnetic field). During the experiment, plasma flow is accelerated outwards from the wire array and expands radially encountering the planar grid. As observed in previous experiments, aluminium wires generate plasma outflows that move at supersonic speed, carrying a fraction of the transverse magnetic field produced during the discharge of the generator current. The collision of plasma with the grid generates a system of reflected shocks in the post-grid region. These shocks have the same spatial periodicity as the grid and the copper rods do not interfere significantly with the magnetic field, as shown in previous experiments using field-aligned obstacles[60].

An overview of the shock dynamics in this experiment can be observed from optical self-emission images taken using a fast frame camera positioned along the same line-on sight of side-on shadowg-

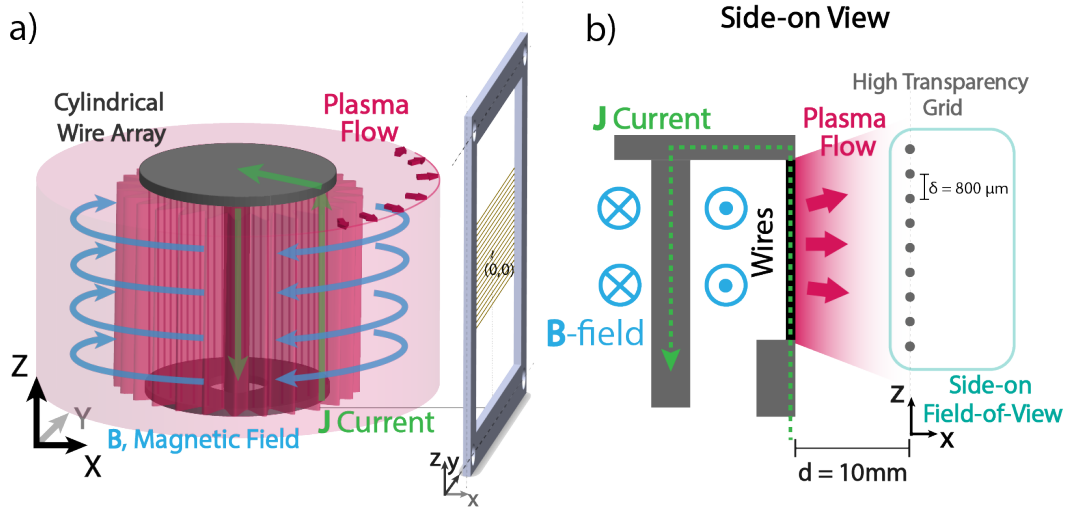


Figure 3.6: (a) Illustration of the exploding wire array with high transparency planar grid perpendicular to the plasma flow direction placed at 10 mm from the wire array. (b) Side-on diagnostic projected view.

raphy and imaging refractometer diagnostics. The fast-framing camera used at MAGPIE is able to capture a sequence of 12 images with variable inter-frame and a minimum exposure time of 5 ns. Further details on this technique will be presented in the next chapters. Examples of self-emission images are shown in Fig. 3.7. The experiment aims to investigate the response of the imaging refractometer technique to a similar transverse density gradient distribution as presented in the synthetic test case above. Therefore, the main interest is in the 2-D plasma structures dominating the side-on view (xz -plane) of the post-grid region. The first frame on the left is a background image taken before the experiment, in the absence of any plasmas. A metallic object was placed in the field of view at a distance of 4 mm from the planar grid as a spatial scale reference. The planar grid is on the left-hand side of the frame at the origin of the horizontal x -axis. The cylindrical wire array is out of the field of view with the shocked plasma starting from the grid and propagating to the right side of the images. At 320 ns from the current start, shocked plasma emerges from the gaps of the grid, forming a system of reflected shocks which intersect with each other along the z -axis as they expand in the vacuum region. In the third self-emission frame taken 80 ns later (400 ns after the current pulse), the shocks appear as elongated plasma structures along the direction of expansion with visible crossing points along the vertical axis. This system of reflected shocks resembles the sinusoidal density distribution presented in Fig. 3.3, making this late-time plasma system an ideal candidate for testing the operation of the imaging refractometer.

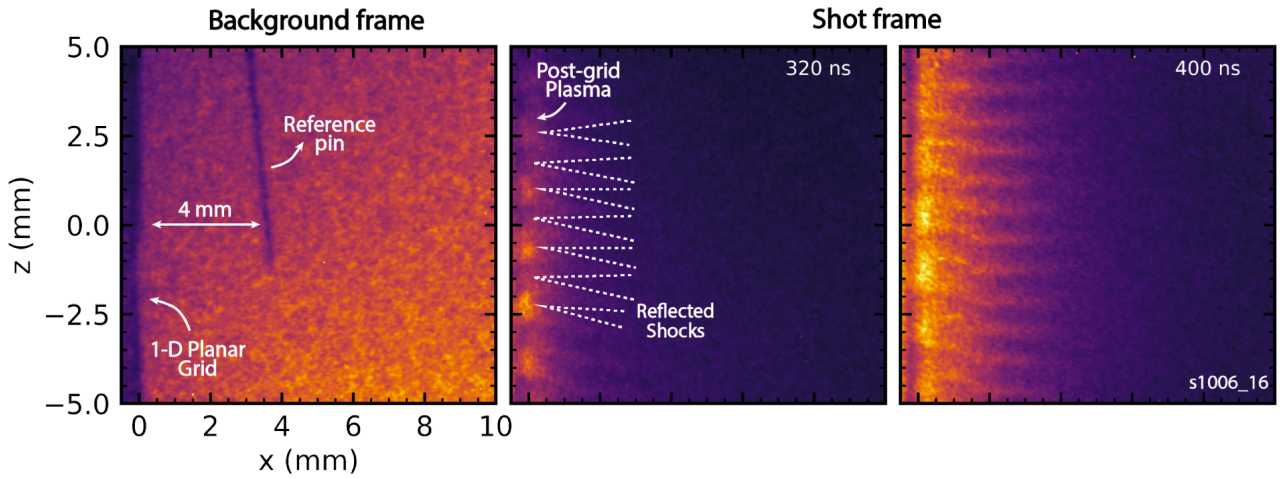


Figure 3.7: Three examples of optical self-emission images taken during the high transparency grid experiment (s1006_16). The first frame is a background image captured before the experiment, in the absence of any plasmas, using an LED lamp as a back lighter. The exposure time is 5 ns and 40 ns inter-frame starting at $t = 160$ ns after the current start. Shown in false colour.

3.4.2 Shadowgraphy and Refractogram Analysis

In this experiment, both shadowgraphy and imaging refractometer diagnostics have been used to image the reflected shock system in the post-grid region at 400 ns after the current start. As observed from the self-emission images, Fig. 3.7, at this late time, the shock structure is clearly visible extending until $x = 4$ mm from the planar grid, showing enough plasma structures to test the response of the refractometer. As the outcome of the imaging refractometer is not intuitive to interpret, it is worth analysing first the correspondent shadowgraphy data shown in Fig. 3.8.

Fig. 3.8(Left) shows a background image taken prior to the experiment without any plasma. The planar grid is visible at the origin of the horizontal axis, which separates the image field of view in the pre-grid region ($x < 0$ mm) and post-grid region ($x > 0$ mm). A metallic obstacle was positioned at 4 mm from the grid as a scaling reference. The shadowgram captured at 400 ns after the current pulse during the experiment is shown in Fig. 3.8(Right). The plasma flow from the wire array propagates from left to right, interacting with the grid at $x = 0$ mm and forming a dense region of accumulated plasma around the obstacle edges. In the post-grid region, for $x > 0$ mm, multiple oblique shocks emerge at $x = 0.5$ mm expanding as the flow propagates to the right. The oblique shocks occupy the entire height of the grid and are evenly spaced by $\delta = 0.8$ mm in the z -direction, corresponding to the grid rod spacing, as it can also be seen from the vertical line-out, labelled as V0. The presence of these

shocks generates strong density gradients which deflect the probing laser beam. Further from the grid, the strength of the shocks reduces due to the combined effect of the decreasing overall plasma density imposed by the current history and the radial divergence of the flows. The intersection points of these shocks are visible in the shadowgraphy image at $x = 0.5$ mm, $x = 1.8$ mm and $x = 4$ mm. An additional interaction point highlighted by the projected lines occurs at $x = 6.7$ mm. This is not directly visible from the shadowgraphy, however, it appears clearly in the refractogram as a broadening effect in the deflection angle distribution, as will be discussed later.

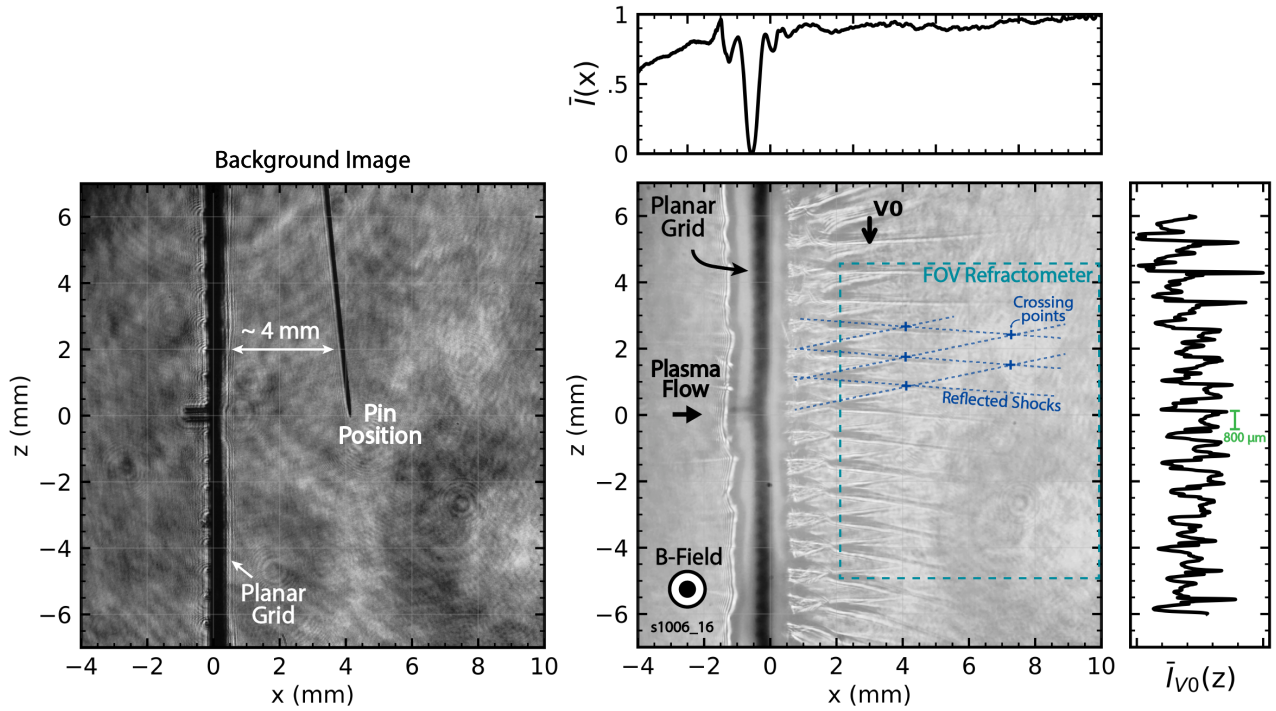


Figure 3.8: 1053 nm side-on laser probe shadowgram of the post-grid region at 400 ns after the starting of the current pulse. The origin along x is set at the 1D grid position with the exploding wire array on the left side of the obstacle (negative values). The zero on the vertical axis corresponds to the half height of the wire array. Multiple crossing shocks are formed along the Z -axis after the passage of the plasma flow through the grid. The intensity profile along the x -direction remains almost constant after the grid as shown from the averaged intensity $\bar{I}(x)$. The intensity distribution in the z -direction from the line-out $V0$, $\bar{I}_{V0}(x)$, is periodic as the imposed planar grid.

The signal of the imaging refractometer captured together with the discussed shadowgraphy data is shown in Fig. 3.9. The image represents spatially resolved measurements of deflections of the probing laser beam. The distribution of deflection angles is symmetric about $\phi = 0$ mrad, as expected by the symmetry of the obstacle. Near the obstacle, at $x = 2$ mm, the envelope of the signal begins very broad

(up to 30 mrad) and decreases as the flows propagate further away from the grid. This is consistent with the weakening of the density gradients. At $x > 10$ mm, the signal width is only slightly broader than the characteristic response of the diagnostic measured in the absence of any plasma prior to the experiment. An additional feature is the presence of strong periodic intensity modulations extended along the horizontal (spatial) axis of the image. These intensity modulations are caused by constructive and destructive interference effects which will be discussed later below. As shown in the previous

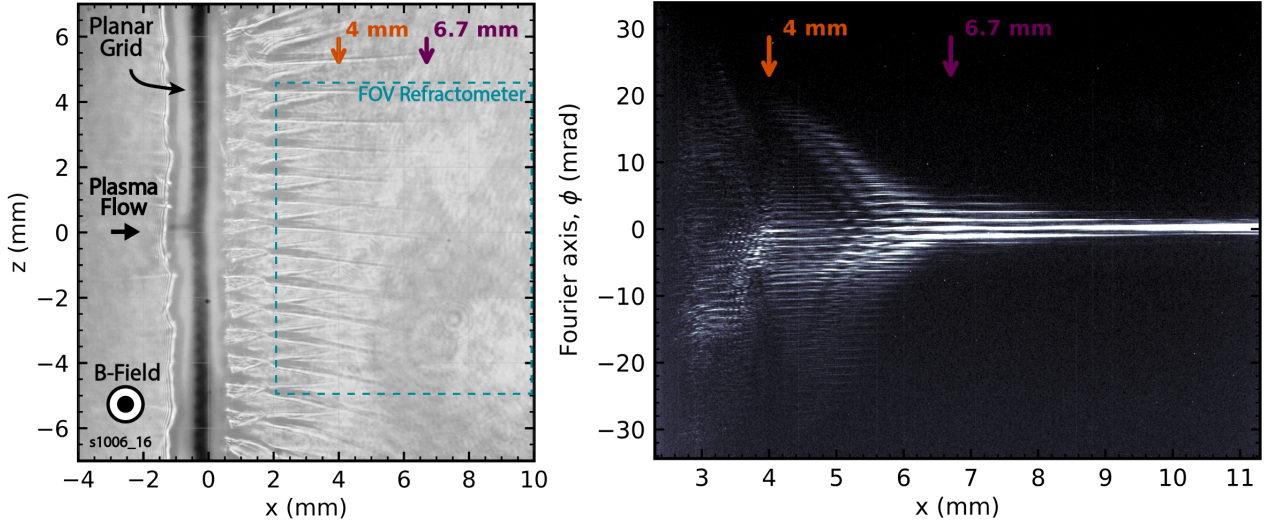


Figure 3.9: (Left) shadowgraphy image presented in Fig. 3.8. (Right) refractography image of the highlighted Field-of-View (FOV) region in the shadowgram. The laser intensity does not saturate the camera detector at any points in the raw image.

section, according to eq. 2.16 stronger density gradients produce larger ray deflections which result in broadening effects of the intensity distribution around its central value. Without any assumptions on the type of distribution, the standard deviation of the signal intensity represents a statistical description of angular deflections for a given spatial position, which can be defined as:

$$\sigma_{\phi} = \sqrt{\sum_i \phi_i^2 I(\phi_i) / \sum_i I(\phi_i)} \quad (3.12)$$

where $I(\phi_i)$ is the intensity of the pixel corresponding to the deflection angle ϕ_i . Experimental refractograms are usually affected by background noise which can strongly affect the estimate σ_{ϕ} if not removed from the raw signal. Since the background is almost constant, this is achieved simply by subtracting the signal baseline from the data. This preliminary signal processing is illustrated in Fig. 3.10. Fig. 3.10(a) shows an example of vertical line-out taken from the refractogram of Fig. 3.9(Right),

representing the raw distribution of the laser deflection angles. The red line represents the baseline of the raw signal which is then subtracted to produce the final deflection angle distribution plotted in Fig. 3.10(b) together with the cumulative intensity distribution of the corrected and normalised signal, Fig. 3.10(c). Two examples of intensity distributions of angular deflections are shown in Fig. 3.11 (a)

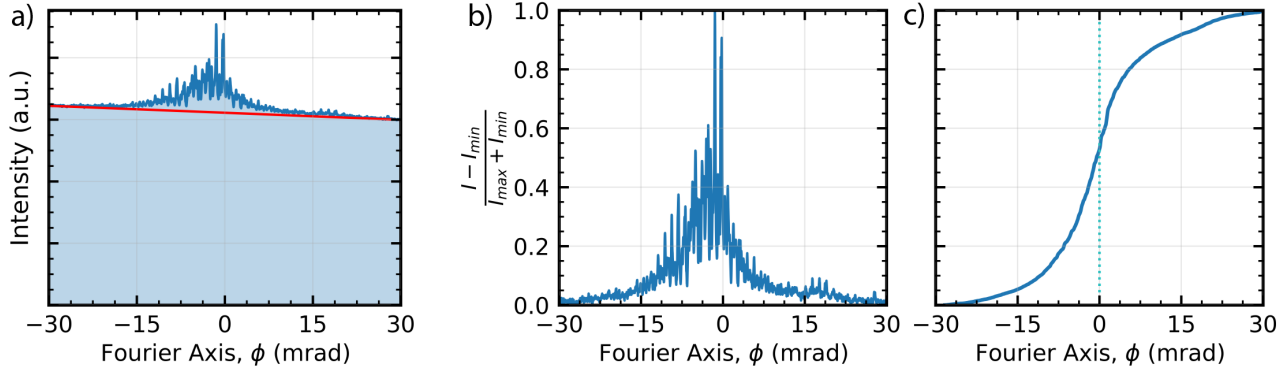


Figure 3.10: (a) Example of raw deflection angle distribution. The red line represents the baseline of the signal. (b) Final deflection angle distribution obtained by subtracting the baseline value from the raw signal. (c) The normalised cumulative intensity distribution of the corrected signal.

& (b) of the image in Fig. 3.9(Right) at spatial position $x = 4$ mm and $x = 6.7$ mm respectively. In both cases, the intensity signal is much broader than the characteristic response of the refractometer, $\sigma_\phi \simeq 0.12$ mrad, represented by the green profile in Fig. 3.11. The recorded broadening of the angular distribution for $x = 4$ mm is $2\sigma_\phi \simeq 7.5$ mrad while reducing at $x = 6.7$ mm to a value of $2\sigma_\phi \simeq 2.5$ mrad. By extending the standard deviation analysis to measure angular deflection distributions at each spatial position x , the overall standard deviation profile can be reconstructed as shown in Fig. 3.11(c). The locations of the oblique shock interactions are visible in Fig. 3.11(c) as a sudden increase in σ_ϕ at $x = 4$ mm and $x = 6.7$ mm, corresponding to the locations seen in the shadowgraphy image. It is notable that the peaks in σ_ϕ occur after the corresponding crossing points measured from the shadowgraphy image. This is because the locations of the shock intersections in the x -axis vary along the z -axis, consistent with the axial expansion of the flows into the vacuum regions above and below the arrays. It is important to note that, although the third shock crossing is not visible in the shadowgraphy image, it can be seen in imaging refractometer data at $x = 6.7$ mm, demonstrating the sensitivity of this diagnostic to very small deflection angles. The broadening profile reconstructed from the analysis shows a similar decay profile to the synthetic sinusoidal perturbation in the previous section (Fig. 3.3), indicating that the capability of the imaging refractometer quantitatively characterises the distribution

of density gradients in a plasma.

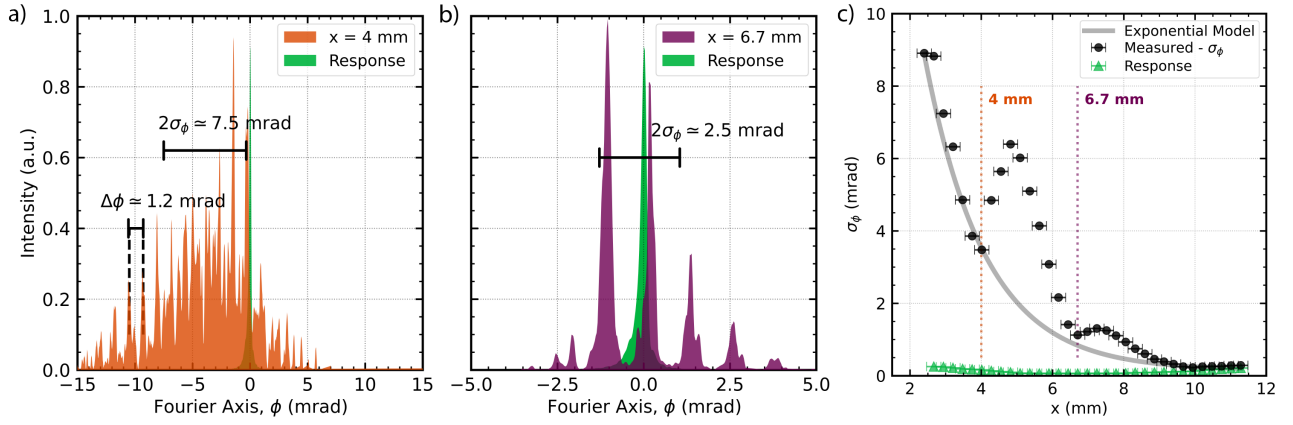


Figure 3.11: (a) and (b) Examples of angular intensity distribution respectively at $x = 4$ mm and $x = 6.7$ mm. $\Delta\phi \sim 1.2$ mrad. Broader angular deflections are observed in the regions closer to the obstacle. Peak displacements resulting from interference effects are spaced by $= 1.2$ mrad, visible in almost the entire refractogram. (c) The measured standard deviation of the refractogram signal, σ_ϕ , at each discrete position with $\Delta x = 0.4$ mm.

A notable feature in the imaging refractometer data presented above is the presence of strong modulations in the distribution of deflection angles (see Fig. 3.9), which can be attributed to axially periodic perturbations in the electron density by the grid. The imaging refractometer data show intensity modulations spaced by $\Delta\phi \approx 1.2$ mrad along the Fourier axis [Fig. 3.11(a)] resembling diffraction patterns caused by periodic opaque structures or diffraction gratings. In this experiment, there is no opaque structure present, so this analogy is somewhat misleading. Instead, each shock acts as a refracting object, deflecting light rays at various angles. The periodicity of the shocks combined with the varying path differences caused by different refraction angles leads to constructive and destructive interference which is then captured by the imaging refractometer camera. The angular periodicity of the interference pattern can be related to the period of the shock system through [97]

$$\Delta\phi_n = \frac{n\lambda_0}{\delta} = n \frac{1.053 \mu m}{800 \mu m} \approx n \cdot 1.3 \text{ mrad} \quad (n \in \bar{\mathbb{Z}}) \quad (3.13)$$

By substituting the known shock period along z imposed by the grid, δ , and the laser wavelength, λ_0 , we find a remarkable agreement between this simple theory and the measured $\Delta\phi \approx 1.2$ mrad. Although a direct relation between this interference effect and the spectrum of the plasma electron density distribution is still missing, the presence of coherent effects suggests uniform density gradients

along the probing beam (y-axis), which is to be expected given that the obstacle is extended in this direction.

3.4.3 Summary

In this chapter, a novel refractive-based diagnostic, called imaging refractometer has been introduced together with a newly developed custom ray tracing tool for generating synthetic shadowgraphy and refractometer images. Experimental validation of the diagnostic was performed using an aluminium exploding wire array and a shock system produced by a planar high transparency grid. The interaction of the supersonic outflow from the wire array with the planar grid at a fixed distance from the wire array forms multiple reflected shocks, expanding in the post-grid vacuum region. Both shadowgraphy and imaging refractometer diagnostics are then used to image this shock system from the side-on view. From the analysis of the raw refractogram, significant broadening was recorded near the grid while decreasing along the spatial axis, as expected from the reduced density gradients due to the combined effect of the decreasing overall plasma density imposed by the current history and the radial divergence of the flows. Quantitative estimates of laser deflection angles have been successfully obtained from the imaging refractometer data, demonstrating the superior sensitivity of the diagnostic to variation in density gradients. Finally, strong modulations in the distribution of deflection angles due to coherent effects, suggest the presence of uniform density gradient along the beam direction with angular periodicity consistent with the modulation imposed by the grid rods. In the next chapter, the imaging refractometer is applied to investigate the structure of reverse shocks formed from the collision of super-magnetosonic plasma against conductive planar obstacles, providing valuable quantitative insights into density structures and the overall underlying physics.

Chapter 4

Instabilities and Structures in Accretion

Shock Experiments

In recent years, significant advancements have been made in our understanding of magnetised shocks in plasmas through laboratory experiments. Various experimental platforms have been used to study these shocks, including laser-driven magnetised collisionless shocks [98], laser-driven plasma flows interacting with magnetised obstacles [99], supercritical perpendicular shocks in theta-pinch configurations [100], and pulsed-power-driven wire-array configurations [90, 101]. Among these platforms, the use of pulsed-power-driven inverse wire arrays has proven to be particularly valuable in investigating shock physics, generating long-lasting, quasi-one-dimensional, supersonic, and often super-Alfvénic magnetised plasma flows. These flows possess characteristics that are scalable to many astrophysical phenomena, including astrophysical jets, accretion shocks and interactions in the solar wind, offering a versatile environment to explore different plasma regimes and fundamental properties of high energy density plasmas.

When studying astrophysical shocks, thermal instabilities and radiative cooling effects can play crucial roles in the evolution of shock systems and their stability, giving rise to oscillating shock fronts and possibly leading to turbulence. This behaviour is especially relevant in highly magnetised, low- β accretion flows onto Young Stellar Objects, where confining magnetic fields effectively make the flows one-dimensional. Recent analyses and simulations conducted in two or three dimensions

have shown that the shock structure in these systems will fragment into multiple oscillating regions, rather than exhibiting quasi-one-dimensional oscillations [26, 27, 58]. Furthermore, the strength and orientation of the magnetic field relative to the accretion flow have been found to influence the stability of the system with respect to shock perturbations. The research presented in this chapter builds upon previous investigations that utilised exploding aluminium wire arrays to explore various aspects of accretion shock formation. These earlier studies examined reverse shocks formed using planar obstacles [59], shocks forming around magnetised obstacles [90], subcritical shocks around large, field-aligned obstacles [63, 65], and bow shocks forming around small, pointed obstacles [60]. The insights and findings from these earlier experiments serve as a foundation for the current set of experiments, which aims to delve into the complex dynamics of magnetised accretion shocks and the role of radiative cooling instabilities in different plasma regimes (specifically using tungsten-based plasma flows to increase the strength of radiation effects of the plasma).

This chapter is structured as follows. We begin by introducing the experimental platform employed in our investigations, along with the diagnostic setup utilised for data collection. Subsequently, we present a summary of the shock dynamics observed in aluminium plasmas, laying the foundation for understanding reverse shocks in high-Z plasmas. The focus is then moved to experiments with tungsten plasma flows, providing firstly a discussion of the reverse shock and secondly, a detailed analysis of the observed instabilities. Lastly, we discuss the density features observed and offer interpretations to the underlying physics behind these phenomena.

4.1 Experimental Setup and Diagnostics

The experimental setup utilised in this section employs a cylindrical wire array configuration driven by the MAGPIE Pulsed Power Generator, as discussed in Chapter 2, combined with flat conducting obstacles. This configuration has proven effective in generating long-lasting plasma flows that persist for approximately 500 ns while advecting a portion of the driving magnetic field. Figure 4.1 provides a diagram of the experimental setup employed in this study. The cylindrical wire array consists of 20 parallel thin metallic wires, connected between two electrodes. The wires are arranged cylindrically,

with the cathode extending along the central axis of the array. The dimensions of the array are fixed at 20 mm in diameter and 22 mm in height, with the obstacles positioned 10 mm away from the wires. To ensure improved flow uniformity toward the target and minimise divergence in the region of interest, the 9 wires facing the obstacles are positioned closer together, with an angular separation of 11.25° , while the remaining wires are spaced at intervals of 22.5° . The wires are suspended from a wire holder, passing through both the cathode disk and the anode ring grooves. They are securely attached to the outer cylinder supporting the anode ring and connected to the inner anode. The entire array is coaxial, with the internal cathode post directly linked to the central conductor of the current feed using a hollow support. This configuration offers an inverse current path compared to a conventional z-pinch wire array, enabling outward acceleration of the plasma.

Access to the end-on laser probe into the imaging optics is facilitated by holes through the anode plates, the cathode and the central support, via a 45° mirror. The inner anode is mounted on the outer anode plate, which is then grounded to the chamber breadboard via 8 stainless steel return posts.

The selection of wire material allows for the exploration of various plasma regimes, enabling changes in ionization state ($\bar{Z}$), resistivity, and radiative cooling while keeping the overall flow parameters such as mass density, temperature, and flow velocity relatively constant. For this series of experiments, the focus is on aluminium and tungsten wire materials, while intermediate-Z elements such as copper and silver are considered potential subjects for future research. The passage of current through the wires leads to the resistive heating of the wire surfaces, resulting in the formation of a coronal plasma while the wire cores remain cold and dense [92]. The current flowing through the central electrode (cathode) generates a strong azimuthal magnetic field that accelerates the plasma flow outward, with a fraction of the generated magnetic field advected by the plasma flow. The ablated plasma flow is characterized by its super-sonic and super-Alfvénic nature, expanding outward in the surrounding region throughout the duration of the current pulse due to the $\mathbf{J} \times \mathbf{B}$ force. For these experiments, a planar obstacle is positioned 10 mm away from the front wires to establish a quasi-one-dimensional flow interaction. The absence of current return structures surrounding the wire array provides convenient access from both side-on and end-on (or axial) views, facilitating the implementation of various spatially and temporally resolved diagnostics. A schematic representation of the experimental setup is shown in Fig. 4.2. The diagnostic setup includes optical self-emission

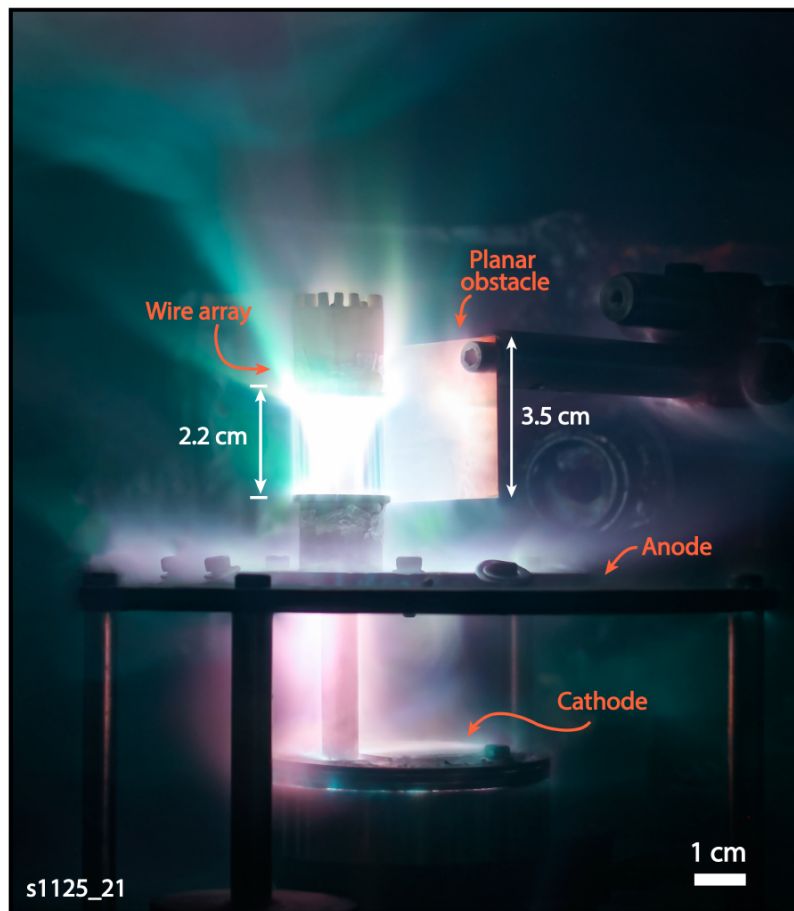
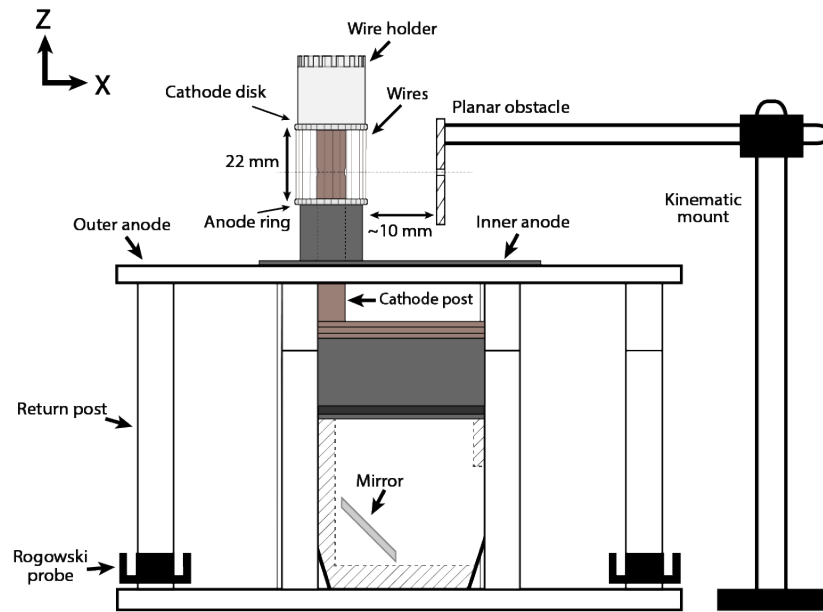


Figure 4.1: (Top) Illustration of the experimental setup. (Bottom) Time-integrated long exposure image taken during an accretion shock experiment at MAGPIE pulse power generator, s1125_21. The image was taken using a simple DSLR Canon 500D with a neutral density filter in front of the optical lens. The shutter is triggered before the pulsed power to avoid any EMPs and kept open for the entire duration of the experiment. Camera settings: ISO-100, 31 mm focal length, 1.3 sec exposure time and f-stop f/5.6.

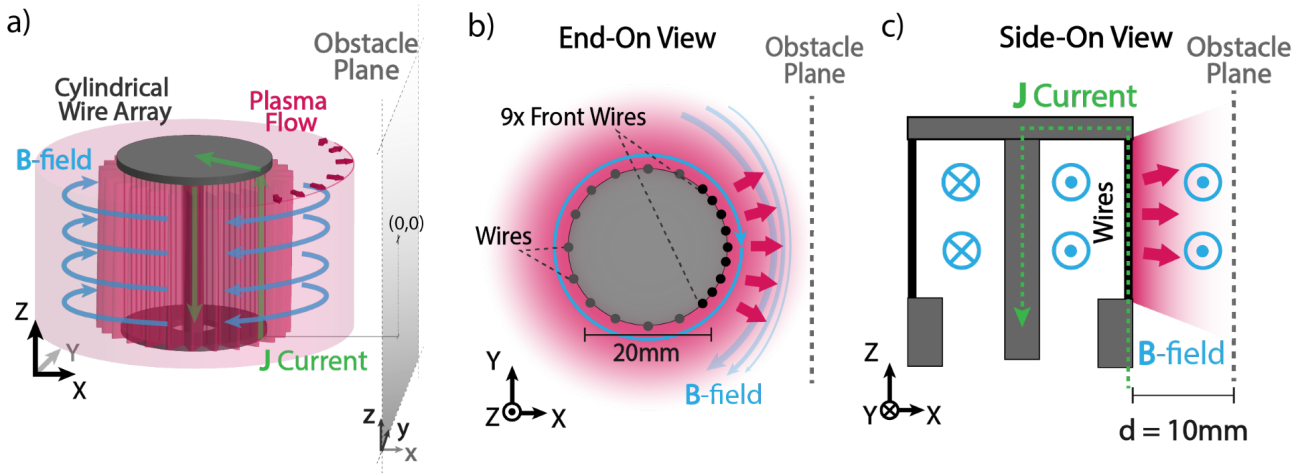


Figure 4.2: (a) Illustration of the exploding wire array experiment with planar obstacles. Obstacles are placed at a fixed distance from the plasma source, on the yz -plane perpendicular to the flow direction. (b) and (c) are the diagnostics projected views, respectively from the end-on and side-on.

imaging from a side-on view and laser-probed interferometry imaging in both the xy -plane and xz -plane, utilising two different wavelengths, 532 nm and 355 nm. Additionally, Thomson scattering diagnostics are employed to further characterise the plasma, specifically the tungsten-based plasma, on the transverse xy -plane at the mid-array height. The complete Thomson scattering geometry is discussed later in this section; however, to accommodate the laser beam along the plasma flow direction, two holes were introduced in the planar obstacle and the cathode post, respectively. In addition to the above diagnostics, conventional machine diagnostics were used to evaluate the performance of the MAGPIE generator. Two pairs of Rogowski probes were positioned in two opposing return posts, as depicted in Fig. 4.1. Each pair of probes was mounted with opposite polarity, allowing the electrostatic pickup from each post to be removed and facilitating two independent measurements of the load current. Figure 4.3 illustrates the time-resolved voltage signals and the corresponding current profiles for the aluminium and tungsten shots discussed in this chapter. The recorded voltage signals indicate the rise in voltage detected by the Rogowski probes, which serve as the reference for determining the starting time of the current pulse in each experiment as well as the reference timing of all the diagnostics presented in this study. By integrating the voltage signals over time and applying a calibration factor, the current profiles (represented by solid-coloured lines) are reconstructed. Additionally, the measurement of the current profile by the Rogowski probes enables a comparison of the drive strength across different shots.

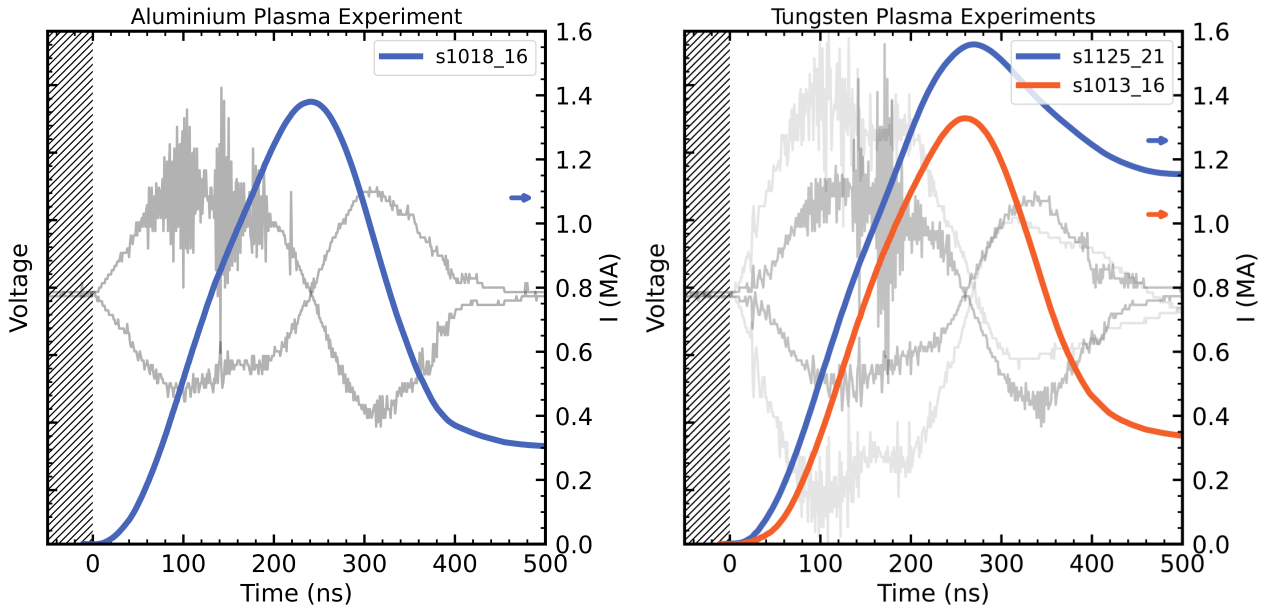


Figure 4.3: Example of current profiles of the experimental shots, s1018_16 for aluminium plasma and s1013_16, s1125_21 for tungsten array measured using two pairs of Rogowski coils. The shaded grey lines represent the raw voltage signals of Rogowski probes recorded during the experiments.

4.2 Reverse Shock Structure in Magnetised Aluminium Plasma

In this section, the reverse shock structure formed from the interaction of supersonic and super-Alvénic aluminium plasma flows against a planar conductive obstacle is presented. This reverse shock system has been studied extensively in previous work, specifically using Thomson scattering measurements and laser probe interferometry [59, 60, 86]. Therefore, only self-emission images and the novel imaging refractometer data will be discussed here, highlighting the main characteristics of reverse shocks produced with this platform.

4.2.1 Overview of the reverse shock dynamic from optical self-emission images

Optical self-emission imaging is a valuable diagnostic tool for investigating the overall dynamics of plasma systems, as it enables multiple images to be captured in a single shot using a high frame rate. This diagnostic technique is straightforward to implement, requiring only a camera and imaging optics. Although quantitative analysis of self-emission brightness in these images is challenging due to its dependence on various plasma parameters, a qualitative assessment can be made by comparing the self-emission images with interferometry data. The correlation between areas of increased brightness

in self-emission images and higher electron density becomes apparent when comparing them with interferometry results. This correlation arises from the quadratic relationship between self-emission intensity and electron density, governed by line emission in optically thin regimes. Consequently, self-emission images offer a reasonable qualitative estimation of density structure evolution during a shot. They also provide guidance for determining when to employ other quantitative diagnostic techniques in future experiments. The combination of qualitative information from self-emission images with quantitative data from other diagnostics enables deeper insights into the underlying physics of the plasma system.

The fast-framing camera used at MAGPIE has the capability to capture a sequence of 12 images with a variable inter-frame and an exposure time of ≥ 5 ns. For these particular experiments, the inter-frame is set to 40 ns, resulting in a total temporal range of 480 ns. This allows us to observe almost the entire dynamics of the experiment, starting from the plasma interaction with the obstacle at approximately 100 ns and continuing until after the end of the driving current pulse at 600 ns. The aim

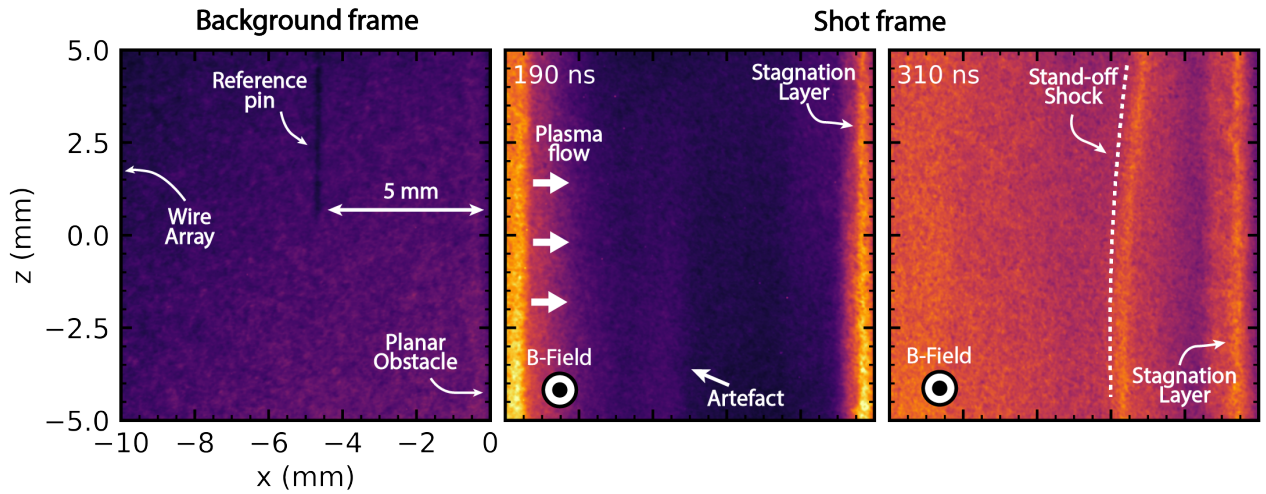


Figure 4.4: Examples of optical self-emission images from the aluminium shot s1018_16. The background frame was taken before the shot using a backlighting LED lamp. The two shot frames are taken at 150 ns and 310 ns after the current start.

of this experimental setup is to investigate the development of a perpendicular shock as a result of the interaction with the planar conductive wall. Therefore, the most informative field of view is from the side-on view in the xz -plane. Although end-on (xy -plane) optical self-emission data are available, we will not discuss them in this context. To ensure a reliable comparison between electron density and self-emission features, the line-of-sight of the camera is aligned with the optical path of the side-on

interferometry. The camera is focused at $y = 0$ mm, and a reflecting bandpass filter is positioned just before the camera, effectively blocking laser light at a wavelength of 532 nm during the experiment.

Typical optical self-emission images for the reverse shock experiments are presented in Fig. 4.4. The first frame on the top left-hand side represents a background image taken before the experiment, without the presence of any plasma. A metallic pin was used as a spatial scale reference placed at a distance of 5 mm from the obstacle surface ($x = 0$ mm). In this case, the illumination was provided by an LED lamp located at one of the side windows of the chamber. In the frame, the wire array is positioned at $x = -10$ mm and during the shot the plasma flows move from left to right, towards the obstacle. The second image represents an early time frame taken at 190 ns after the current start. At this time, the plasma flow produced by the ablating wires from the array expands towards the obstacle, forming a stagnation layer. Part of the magnetic field generated by the inverse wire array is advected by the plasma flow and it piles up in the region ahead of the obstacle surface. This accumulation of magnetic field, with B-field along the y-axis (into the page), perpendicular to the flow propagation, mediates the formation of the stand-off shock visible in the frame of Fig. 4.4(Right) taken at 310 ns after the current start.

The general evolution of the plasma system in a single experiment can be seen in Fig. 4.4, showing a sequence of 12 images captured during the interaction of an aluminium plasma with the planar obstacle. The first frame is taken at 150 ns after the start of the current pulse. The front wires of the exploding array are visible at $x = -10$ mm, with respect to the position of the planar obstacle. At this time, the wires have already started to ablate with the plasma flow starting to accumulate at the obstacle surface, forming a stagnation layer as evident from the image brightness. This dense region expands backwards due to the continuous inflow of material from the array, reaching $x = -0.5$ mm after 80 ns.

At 230 ns after the current start, a new feature emerges at $x = -3$ mm from the obstacle. This marks the early formation of a stand-off shock as a result of the interaction between the plasma flow and the pile-up of the magnetic field advected by the flow. As the flow collides against the target, the magnetic field is unable to diffuse through the obstacle due to the significant difference in spatial scale between the conducting obstacle ($\simeq 35$ mm) and the diffusion length scale of the magnetic field ($\simeq 1$ mm). This stand-off shock remains consistently present throughout the experiment, positioned

s1018_16 - Aluminium Planar Obstacle Experiment - Optical Framing images

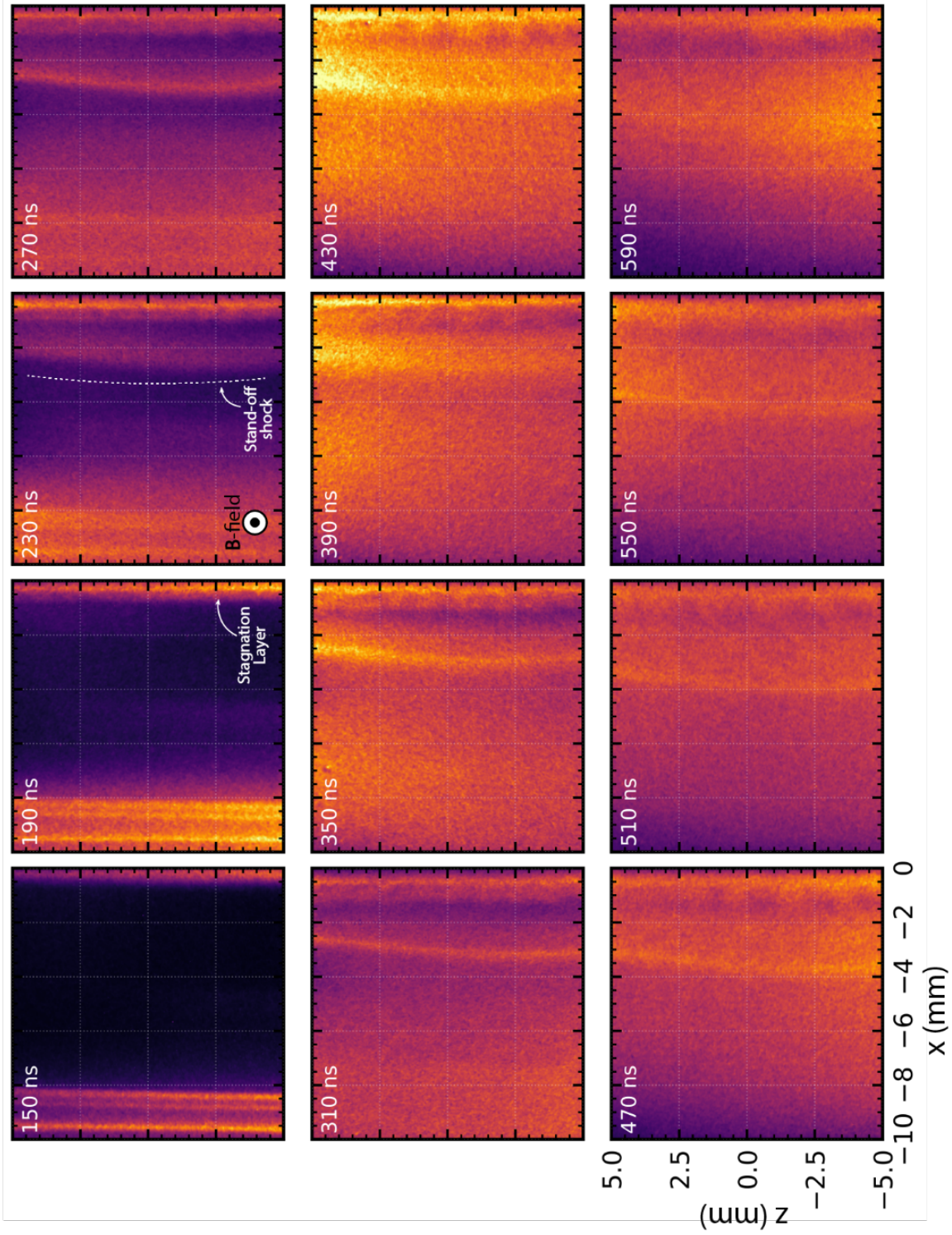


Figure 4.5: Optical fast framing images taken during the shot s1018_16, 5 ns exposure, 40 ns interframe, 12 frames beginning at $t = 150$ ns after the current start. Shown in false colour. Formation of the hydrodynamic shock and development of a stand-off shock mediated by the magnetic field

at approximately $x = -4$ mm from the obstacle, as observed in the remaining self-emission frames (Fig. 4.5). This observation indicates that a pressure balance is established across the shock, where the ram pressure and magnetic pressure are comparable with the thermal pressure playing only a partial contribution. This phenomenon has been extensively investigated in previous research, and for a more in-depth analysis, we refer to the relevant literature [59, 60]. However, an aspect that requires further investigation is the role of the stagnation layer in the overall dynamics of the shock system. Early Thomson scattering measurements performed in [59], showed a small level of velocity and density jump across the stand-off shock at early times, 180 ns and 225 ns after the current pulse. Therefore, it

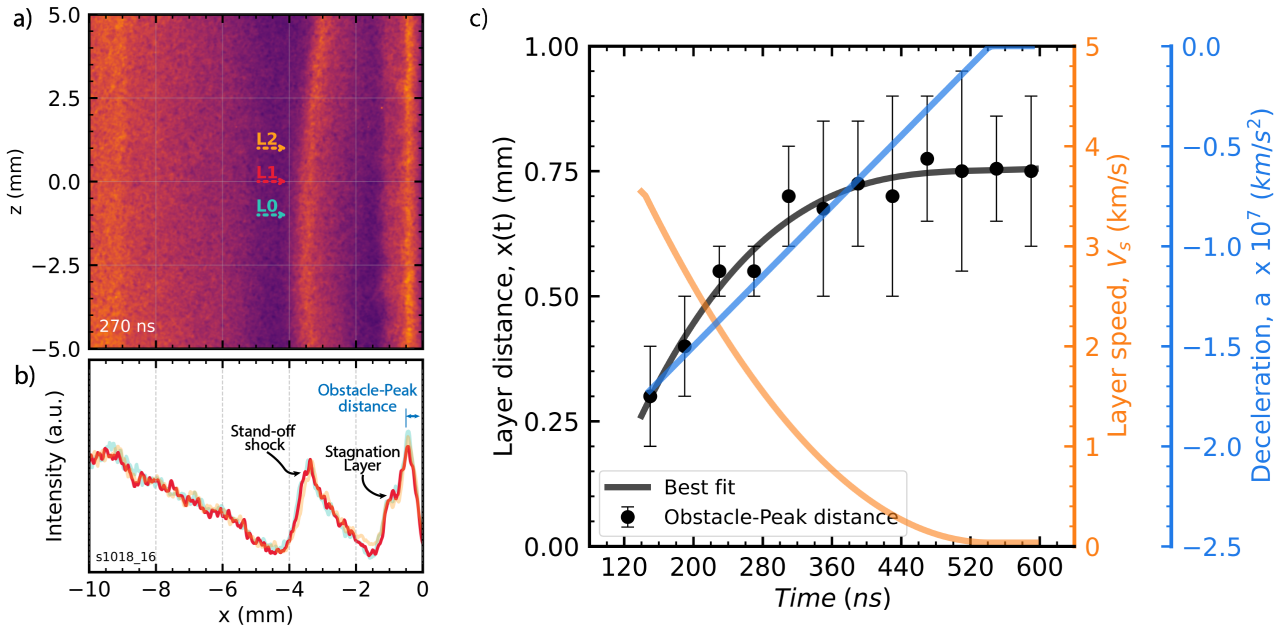


Figure 4.6: (Left) Example of optical self-emission frame taken at 270 ns after the current start. Three horizontal intensity line-outs are used to better identify density structures within the image. (Right) Reconstructed reverse shock dynamic from shock distance measurements estimated as spatial distance between the stagnation intensity peak and the obstacle surface.

is important to note that the observed stagnation layer must be mediated by a hydrodynamic shock formed near the obstacle surface, as the plasma inflow remains supersonic during the early evolution of the system. However, this shock is not resolvable by the current diagnostics capability, so we will refer only to the stagnation layer as a whole. Unlike the stand-off shock, which extends outward from the obstacle surface, the stagnation layer does not exhibit considerable expansion and remains in close proximity to the obstacle wall. Despite the limitations of obtaining quantitative insights from optical self-emission data, it is still possible to estimate the velocity of the stagnation layer expansion by analysing the intensity feature evolving over time. In the example frame of Fig. 4.6(a) taken at

270 ns after the current start, three horizontal line-outs have been taken by averaging the intensity values over 0.5 mm at $z = 1$ mm, $z = 0$ mm and $z = -1$ mm. Intensity distributions are shown in Fig. 4.6(b). From the line-outs, the intensity decreases as the plasma flow expands towards the obstacle, and two prominent peaks representing the stand-off shock and the stagnation layer become apparent as well-defined peaks at around $x = -4$ mm and $x = -0.5$ mm. To estimate the expansion velocity of the stagnant region, the distance between the obstacle and the peak of the stagnation shock in each frame is measured. The resulting measurements are plotted in Fig. 4.6(c). The uncertainties in the data are determined through the assessment of the leading edge of the layer boundary in each frame. To interpolate the intermediate positions, the measured distances are fitted with a cubic model that reaches a plateau at 500 ns. The analysis reveals that the speed decreases, showing a linear deceleration of $\simeq 10^7$ km/s². This non-constant decelerating motion arises from two factors: changes in mass flow due to the finite current pulse rise time, which becomes significant at these timescales, and the presence of a stand-off shock that slows down the plasma before reaching stagnation. Additionally, an interesting observation is the subtle appearance of modulations along the vertical z -axis within the brightest and thinnest region of the stagnation layer. These structures resemble waves that perturb the stagnation layer, as also observed in the shadowgraphy image, Fig. 4.7 and the refractometer image discussed later. Although direct measurements within this dense and cold plasma region are lacking, it is highly probable that these perturbations arise from uneven pressure distribution across the transition, possibly attributable to radiative cooling effects taking place.

4.2.2 Shadowgraphy and refractometer data of aluminium reverse shock

A valuable qualitative diagnostic technique for observing density structures in the presence of plasma density gradients is laser shadowgraphy, as discussed in Chapter 2. The shadowgraphy technique allows imaging redistribution of uniform laser light due to changes in the plasma refractive index. In this section, details of the reverse shock structure observed in the optical self-emission images are presented using data from both shadowgraphy and imaging refractometer techniques taken simultaneously in a single experiment.

An example of a side-on experimental shadowgraphy image is presented in Fig. 4.7, captured

during the aluminium plasma experiment (s1018_16) at 391 ns after the start of the current pulse.

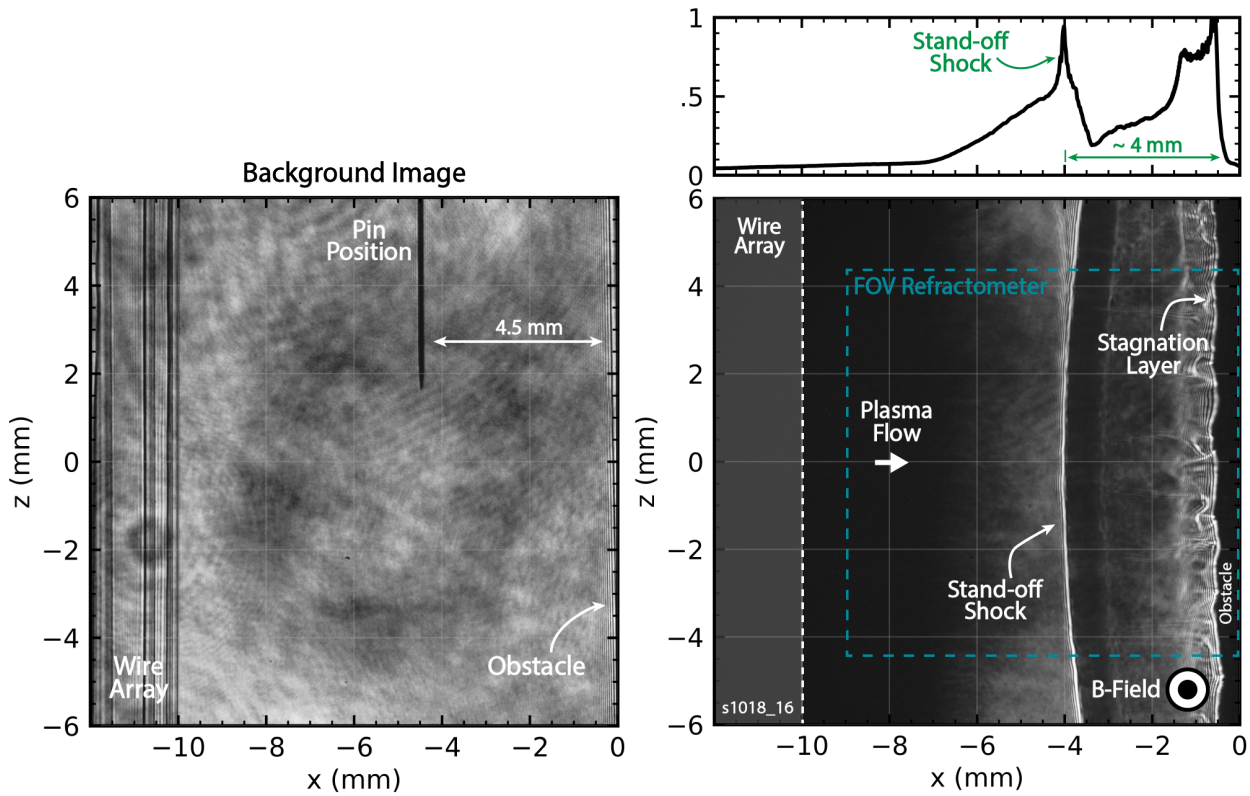


Figure 4.7: (Left) Background image taken before the shot with a pin positioned at 4.5 mm from the planar obstacle at the x-axis origin. (Right) 1053 nm side-on shadowgraphy image of aluminium plasma flow colliding against a planar obstacle, taken at 391 ns after current start. The origin of the horizontal axis is assumed to be at the obstacle position which was placed at 10 mm from the exploding wire array. The stand-off shock is formed at $x = -4$ mm. Visible caustics are present at the beginning of the stagnation layer, at $x = -1$ mm from the planar wall.

Fig. 4.7(Left) shows a background image taken prior to the experiment in the absence of plasma. A small metallic pin was placed as a reference for scaling, positioned at a distance of 4.5 mm from the origin of the horizontal axis corresponding to the location of the obstacle. The wire array is visible on the left side of the image, with the front wires located at $x = -10$ mm from the obstacle surface. Fig. 4.7 (right) shows the shadowgraphy image of the plasma flow interacting with the flat obstacle positioned at the origin of the x-axis, $x = 0$ mm.

Consistent with the optical self-emission images presented in Fig. 4.5, a stand-off shock is observed at $x = -4$ mm, resulting from the accumulation of the magnetic field perpendicular to the flow direction. Additionally, a dense stagnated region is observed near the surface of the obstacle, corresponding to $x > -1$ mm. In the region before the shock, specifically in the range of $-10 \text{ mm} < x < -6 \text{ mm}$,

the shadowgraphy signal experiences significant attenuation. This attenuation is attributed to the strong inverse bremsstrahlung absorption of the probing beam. It is worth noting that even though the plasma density in this region is expected to be below the critical density for the probing wavelength ($n_{cr} = 10^{21} \text{ cm}^{-3}$), the probing beam is still absorbed. As we move away from the array, the plasma density decreases, allowing the beam to pass through unattenuated towards the camera. Notably, a distinct curved bright region is clearly visible around $x = -4 \text{ mm}$ due to shadowgraphy effects from strong density gradients in this region. This brightness corresponds to a significant jump in plasma density, which is attributed to the formation of a stand-off shock. The shock occurs as a result of the accumulation of mass and magnetic field against the conducting planar obstacle, leading to a pronounced increase in plasma density. Next to this bright region, there is a corresponding dark region, from which the laser light has been refracted due to density gradients.

Further downstream, the intensity distribution becomes relatively uniform until the beginning of the stagnation region at $x = -1 \text{ mm}$. In the layer interface region between -1 mm and -0.5 mm , significant perturbations in the laser intensity are observed. These perturbations exhibit caustic structures that are consistent with the presence of large-density gradients. Within the stagnation layer, very close to the obstacle ($x > -0.5 \text{ mm}$), the camera does not capture the laser beam due to strong laser deflection. This is attributed to a combination of refraction and possibly absorption effects caused by the stagnated plasma on the surface of the obstacle. While shadowgraphy images provide valuable qualitative information about plasma structures, performing quantitative analysis is challenging because of the limited dynamic range of the diagnostic.

To overcome this limitation, the novel imaging refractometer technique described in Chapter 3 is applied along the shadowgraphy line to provide more quantitative information. Here, a distribution of deflection angles is instead recorded by imaging the Fourier plane on the vertical axis and the image plane on the horizontal axis of the CCD camera. This results in a spatially resolved intensity distribution of the angular deflections induced by the plasma density gradients transverse to the laser probing path. The refractometry image captured during the same experiment and at the same time as the shadowgraphy image is shown in Fig. 4.8. Similarly to the shadowgraphy image, the refractometry signal is only visible for $x > -7 \text{ mm}$ due to strong absorption. Upstream of the stand-off shock, the refractometer signal displays a relatively narrow distribution of deflection angles that gradually becomes

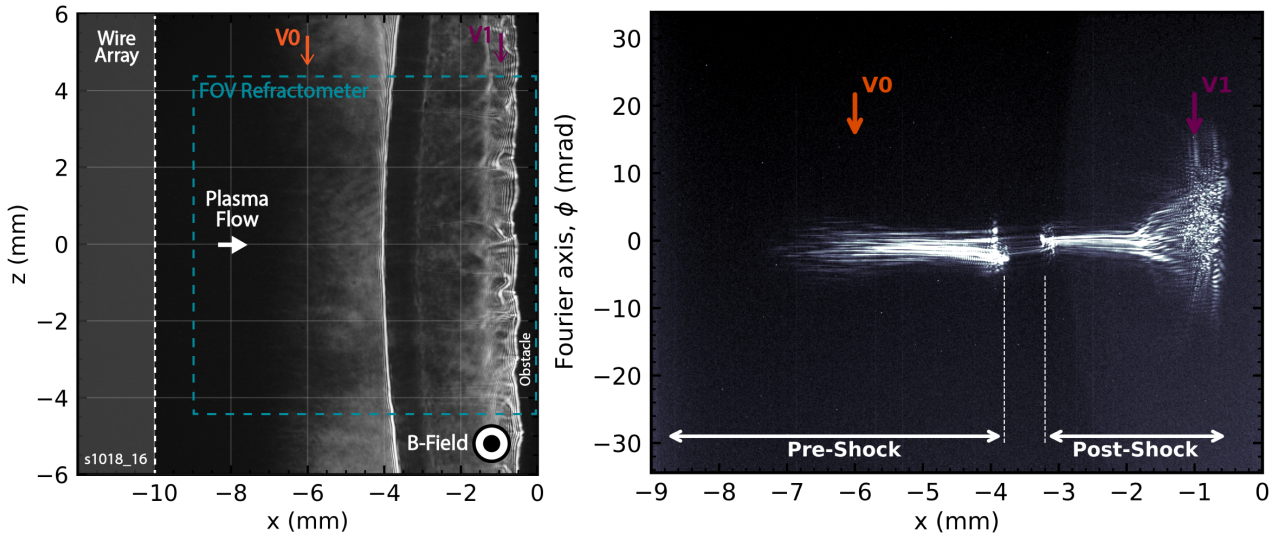


Figure 4.8: Refractogram (right image) of the FOV region highlighted in the shadowgraphy image taken at the same time as the shadowgraphy of Fig. 4.7. Two vertical line-outs labelled V0 and V1 show the angular distributions at spatial positions $x = -6$ mm and $x = -1$ mm in Fig. 4.9. The real dynamic range (for a 16-bit CCD, 2^{16} level of gray-scales, 96 dB) of the data is much larger than can be presented on a printed image so it is better to refer to intensity line-outs.

narrower as we approach the shock though exceeding the angular width of the non-disturbed probe beam. Additionally, the angular intensity distribution exhibits clear modulations along the vertical, angular axis, which arise from the periodic density modulations in the plasma flow, known to occur in the wire array ablation process. These structures and intensity modulation will be discussed later in this section. In the shock transition, the strong density gradient causes a loss of almost all the signal due to shadowgraphy effects within the range of -3.9 mm to -3.2 mm. In the post-shock region, the total signal width becomes narrower, approximately 50% lower, with indications of elongated coherent structures remaining from the upstream perturbation. In close proximity to the obstacle, the distribution of deflection angles becomes much broader and the intensity modulations appear less coherent. Beyond the stagnation interface, specifically at $x > -1$ mm, the overall angular deflection signal reaches its maximum width, while small-scale modulations dominate the intensity distribution. Figure 4.9 showcases two examples of angular deflection distributions obtained from the pre-shock and post-shock regions taken from the vertical line-outs V0 and V1. Both distributions exhibit significantly broader angular spreads compared to the characteristic response of the refractometer (Green profiles, $2\sigma_\phi = 0.01$ mrad). Notably, the width of the deflection angle distribution more than doubles from $2\sigma_\phi \approx 3$ mrad at $x = -6$ mm (upstream of the shock) to $2\sigma_\phi \approx 7$ mrad at $x = -1$ mm (downstream of

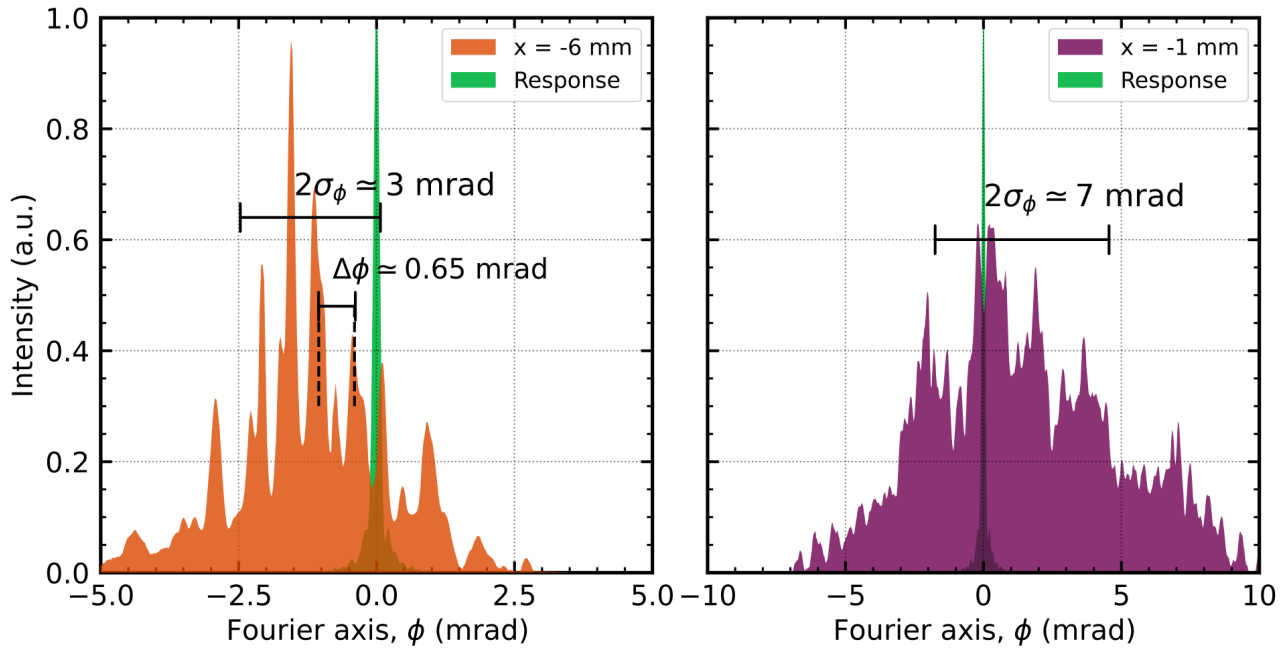


Figure 4.9: Example of angular deflection distribution at $x = -6$ mm (left) and $x = -1$ mm (right), respectively. Broader angular deflections are observed in the regions closer to the obstacle. Peak displacements resulting from interference effects are spaced by $\Delta\phi = 0.65$ mrad, visible both upstream of the stand-off shock and just upstream of the obstacle.

the shock). In addition to the broad angular distribution, it is seen that there is a presence of periodic intensity modulations with deflection angles. These distinct and well-defined peaks in the angular distribution prior to the stand-off shock suggest that there is a periodic density perturbation in the plasma flow. This feature is similar to modulations observed in the planar grid experiment, discussed in Chapter 3, however here there is no grid to impose an axially periodic density perturbation. thus the observed modulation in angular distributions can be attributed to a quasi-periodic modulation of the ablation from wires in a wire array. Previous studies in wire arrays [91, 92, 102], both exploding and imploding configurations, have revealed the rise of shear-flow-stabilised $m = 0$ instability in the produced plasma flows. This instability saturates at a single characteristic wavelength and is strongly dependent on the size of the wire cores, which for low-Z elements, such as aluminium, is observed to be larger than for high-Z materials, like tungsten. Similarly to the case with the grid experiment presented in Chapter 3, section 3.4, the separation of the peaks in Fig. 4.9(a) gives an estimate of the characteristic wavelengths of the density modulation in the upstream plasma flow. Taking a characteristic peak separation of $\Delta\phi \approx 0.65$ mrad as illustrated in Fig. 4.9(Left), we estimate the density modulation wavelength as $\simeq 1.6$ mm using the formula (3.13). Fig. 4.9 shows the presence of various peak

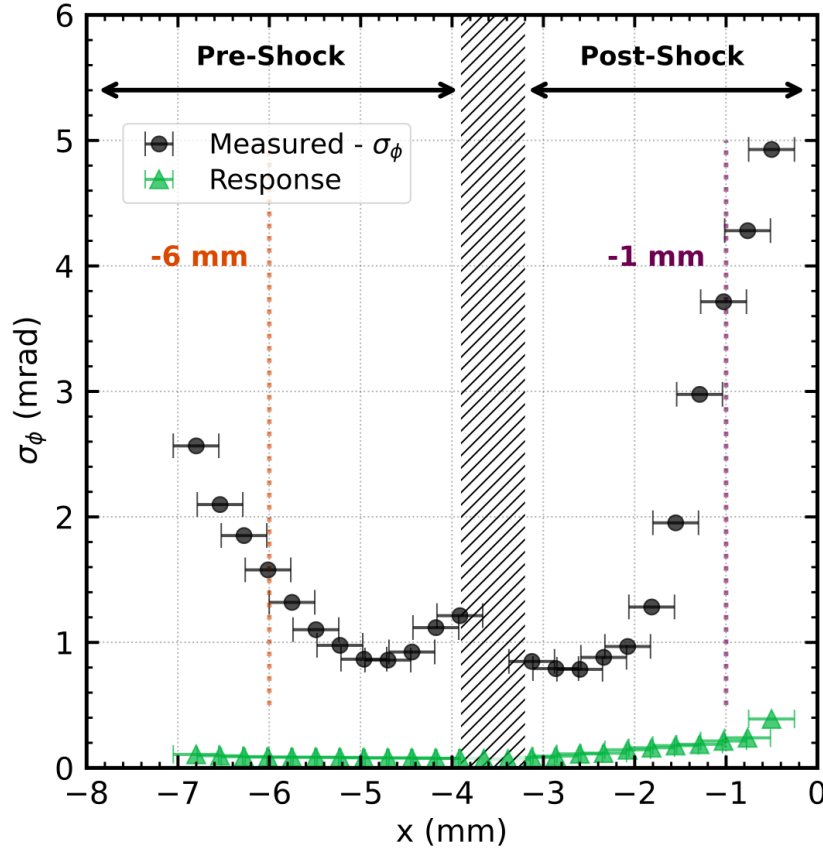


Figure 4.10: The measured standard deviation of the refractometry image signal, σ_ϕ , taking $\Delta x = 0.4$ mm. Almost no signal is detected near the shock transition (masked region).

separations, in the range between $0.5 \text{ mrad} < \Delta\phi < 0.8 \text{ mrad}$, which correspond to the presence of a range of wavelengths in the density modulation. This is natural to expect and consistent with the non-monochromatic density perturbation in the ablation flow [102]. It is worth highlighting that these intensity modulations are not visible in the shadowgraphy diagnostic, but the superior sensitivity of the imaging refractometer reveals these subtle features.

Returning to the observed width of the angular distribution, we extend the analysis of the angular deflection broadening for various spatial positions along the x-axis of the refractogram. Fig. 4.10 provides the standard deviation of the angular deflection distributions, denoted as σ_ϕ , for each spatial position, averaged over spatial intervals of $\Delta x = 0.4$ mm. It is seen that in the upstream plasma region, the width of the angular distribution gradually decreases with x, reaching a minimum of approximately 0.8 mrad at $x = -5$ mm. However, just before the stand-off shock, the angular width increases between $x = -5$ mm and $x = -4$ mm to $\simeq 1.2$ mrad. Immediately after the stand-off shock ($x = -3$ mm $\div$ -2 mm), the angular width drops to 0.8 mrad before starting to increase, reaching $\simeq 5$ mrad in the vicinity

of the obstacle ($x \simeq 0.5$ mm). The implications of these measurements will be discussed later, together with the results from tungsten data.

4.2.3 Summary

In summary, there are several key points to consider regarding the reverse stagnation shock formed in the collision of aluminium plasma flow with the planar conductive obstacle:

- A stand-off shock is formed at $x = -4$ mm, previously interpreted as due to the pile-up of the magnetic field advected by the flow [59]. At the obstacle surface, a stagnation layer is formed, expanding very slowly upstream as material accumulates in the post-shock region. The expansion occurs at a velocity of a few km/s but gradually decelerates over time, eventually reaching a stationary position approximately 0.7 mm away from the obstacle surface.
- The density is smooth in the region just after the stand-off shock and perturbations are observed in the thin stagnation layer ($x = -1$ mm), with wavelength $\lambda \simeq 1$ mm, in both optical self-emission and shadowgraphy images. Direct evidence of the underlying physics causing these perturbations is currently unavailable. However, it is hypothesized that radiative cooling effects within the layer may contribute to their appearance as discussed further in tungsten plasma experiments.
- Refractometer data of the experiment revealed smooth density gradients in expanding plasma flow from the wire and in the region just after the stand-off shock. Coherent interference effects appear in the upstream flow, indicating a periodic mode while intensity fragmentation at the stagnation interface suggests the presence of multi-scale density structures.

4.3 Reverse Shock Structure in Tungsten Plasma

In the previous section, an overview of the experimental setup, diagnostics and reverse shock experiments with an aluminium wire array was presented. The plasma outflow formed through the ablation of thin aluminium wires expands at supersonic and super-Alfvénic speeds, colliding with a flat

obstacle and forming a reverse stagnation shock. The magnetic field carried by the plasma accumulates in front of the conducting obstacle, creating a stand-off shock at a fixed distance from the target surface, supported by the magnetic pressure ahead of the stagnation region. The aim of these experiments is to further explore the features of accretion shocks in the laboratory, specifically investigating the role of the ionisation state and cooling effects on the reverse shock structure observed in the experiments above. To access diverse plasma regimes while maintaining similar hydrodynamic properties such as velocity and mass flow in comparison to previous similar experiments, a tungsten wire array is used as a plasma source instead of aluminium wires.

In this section, a comprehensive analysis of self-emission images obtained from tungsten reverse shock experiments is first presented. We begin by examining the stability of the stagnation layer and highlighting the key differences compared to the previous experiments conducted with aluminium. Second, we proceed to characterize the tungsten plasma flow using both interferometry and Thomson scattering techniques. Additionally, the identification and characterization of small-scale density structures that influence the upstream flow are performed using both shadowgraphy and the newly developed imaging refractometer. Finally, a discussion of the potential mechanism responsible for the observed plasma features is provided.

4.3.1 Overview of the reverse shock dynamic from optical self-emission images

The analysis of the system dynamic has been conducted using a subset of the available optical self-emission dataset, specifically s1013_16 and s1125_21, chosen from various experiments conducted in the past few years. While similarities can be observed between these two tungsten shots, the overall shock evolution is significantly different from that in aluminium plasma. To ensure a consistent comparison, the same camera settings were used across different shots, with an interframe time of 40 ns and an exposure time of 5 ns. In both the aluminium and tungsten cases, the horizontal axis is aligned with the flat obstacle surface in the xz -plane, where the front wires of the array are located at $x = -10$ mm. The vertical axis is set at half the height of the array.

Typical self-emission images of tungsten plasma interacting with a planar obstacle are shown in

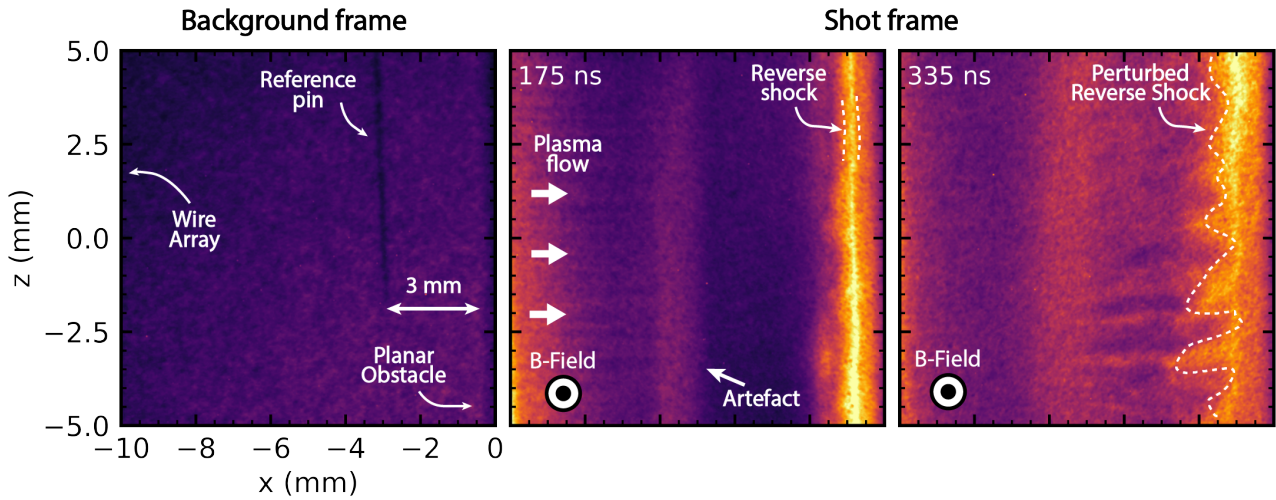


Figure 4.11: Examples of optical self-emission images from the tungsten shot s1013_16. The background frame was taken before the shot using a backlighting LED lamp. The two shot frames are taken at 175 ns and 335 ns after the current start.

Fig. 4.11. Similar to the aluminium experiments, the background image taken before the shot shows the front wires of the array at $x = -10$ mm while the obstacle is placed at the origin of the x -axis. At 175 ns after the current start, a reverse shock structure is observed near the obstacle with a visible high emitting slab at $x = -1$ mm. This bright narrow region extends for the entire height of the obstacle and can be interpreted as the shock transition, separating the upstream flow from the post-shocked plasma near the obstacle. We refer to this post-shocked region as the "stagnation layer" bounded by the brightest narrow slab referred to "thin-shell layer" or "stagnation shock". Behind this narrow slab, a stagnated plasma region is formed at the obstacle surface. Later in time at 335 ns after the current pulse, the reverse shock structure appears modulated along the vertical z -axis with no evidence of a stand-off shock mediated by the transverse magnetic field as it is observed in aluminium plasma experiments. It is important to note that in Fig. 4.11, the bright region at $x = -6$ mm is an unfortunate artefact caused by reflected light from the array, hence does not hold any physical meaning in the data.

The complete sequence of optical self-emission images is shown in Fig. 4.12 and Fig. 4.13 from two different experiments. In the initial frame, which is approximately 140 ns after the current pulse, the plasma flow has already reached the obstacle, and the stagnation region is clearly visible as a bright, narrow region at $x = 0$ mm. Density modulations extending along the x -axis are visible in the upstream flow near the array, up to 255 ns after the current pulse (Fig. 4.12). These perturbations can be attributed to the natural modes imposed in the upstream flow by the exploding wire arrays

s1013_16 - Planar Flat Obstacle - Optical Framing images

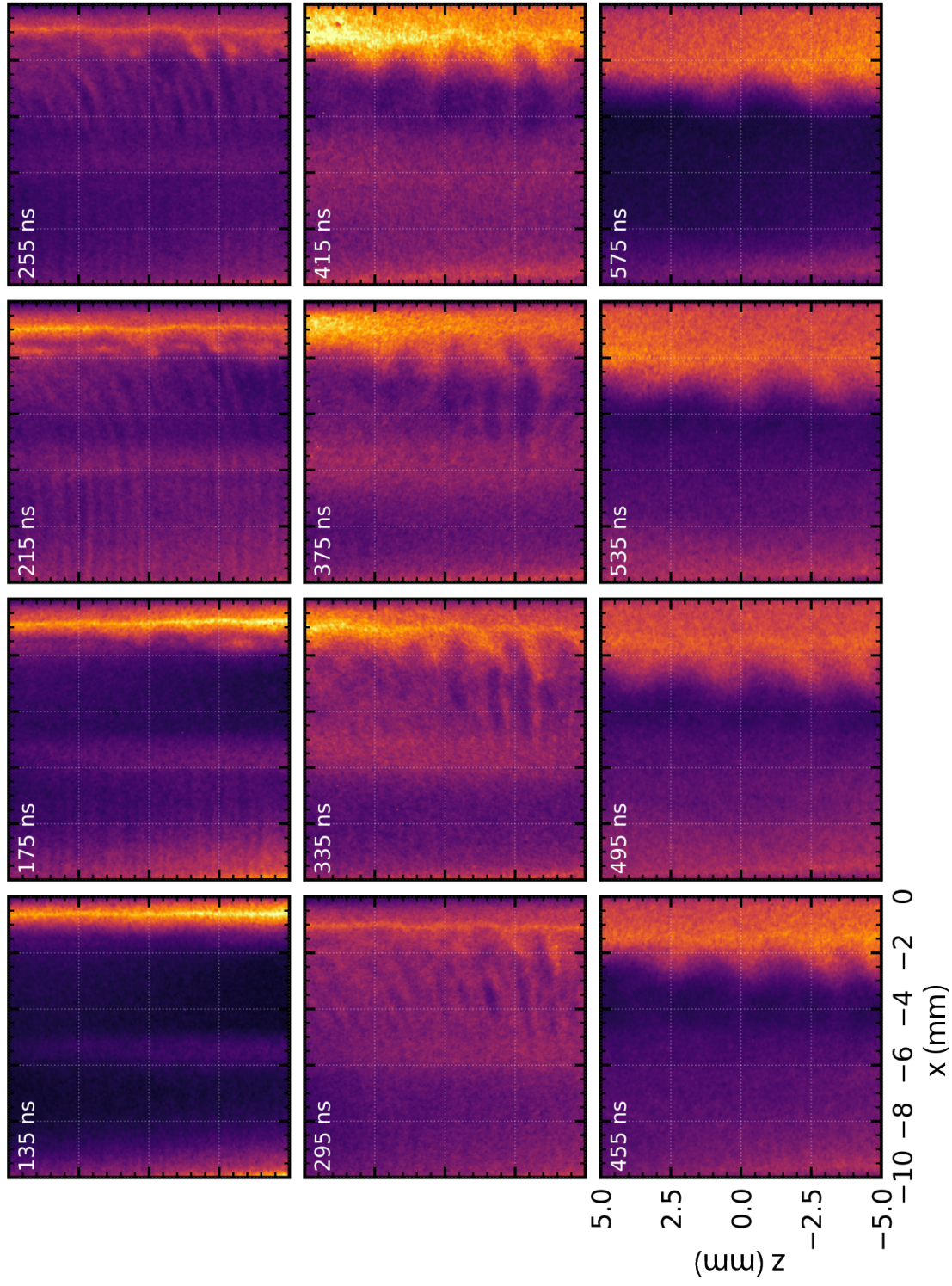


Figure 4.12: Optical fast framing images taken during the shot s1013_16, 5 ns exposure, 40 ns interframe, 12 frames beginning at $t = 160$ ns after the current start. Shown in false colour. Formation of the hydrodynamic shock and development of density structures.

s1125_21 - Planar Flat Obstacle - Optical Framing images

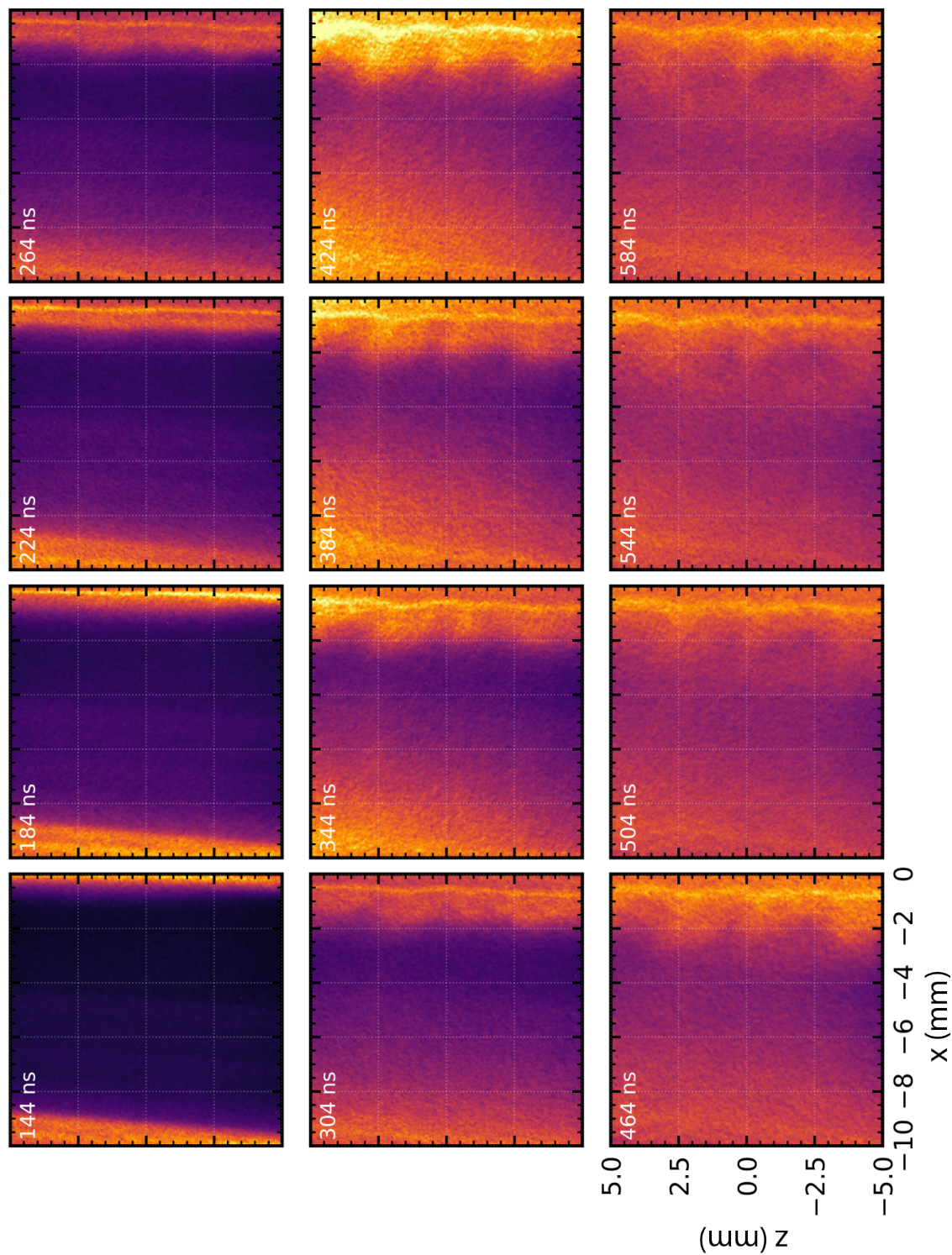


Figure 4.13: Optical fast framing images taken during the shot s1125_21, 5 ns exposure, 40 ns interframe, 12 frames beginning at $t = 176$ ns after the current start. Shown in false colour. Formation of the hydrodynamic shock and development of density structures.

during the ablation phase of the wires. Details on these perturbations will be provided later in the chapter when discussing the shadowgraphy and imaging refractometer data. As time progresses, the stagnation layer expands upstream due to the accumulation of post-shocked plasma near the surface of the obstacle. The stagnation layer propagates upstream at a relatively low speed, reaching a final distance of approximately 1 mm away from the surface obstacle at 300 ns after the current pulse. The propagation speed is comparable to the one observed in aluminium plasma which suggests a similar flow velocity profile. There is no sign of a stand-off shock forming away from the obstacle, instead, perturbations are dominant both at the leading edge of the stagnation layer and in the upstream flow. Looking closely at the interaction region, oscillations appear to corrugate the shock front, increasing in amplitude as the shock expands back to the upstream. This is evident in Fig. 4.12 and even more pronounced in Fig. 4.13. The stagnation shock exhibits a distinct quasi-two-dimensional profile represented by the brightest feature within the interaction region. It spans the entire height of the obstacle and is characterised by an increasing thickness. Similar to the behaviour observed in the aluminium experiment, this thin-shell structure displays corrugations with a clearly defined wavelength. As the stagnation layer expands into the upstream region, these corrugations increase in amplitude until reaching a plateau at 424 ns. Notably, the overall structure of the stagnation shock remains intact throughout the experiment, observable until the last available frame.

Unlike aluminium plasma, the reverse shock in tungsten is directly embedded in the dense plasma region near the obstacle which prevents us from performing any laser-based diagnostic measurements across the shock due to steep density gradients. Nonetheless, quantitative insights can be obtained from the intensity structures in the self-emission images which can help to further support the physical mechanism of the shock dynamics, bearing in mind that both line-integrated effects and the limited resolution of the camera can significantly compromise the accuracy of the results.

Stagnation layer analysis using self-emission images

The collision between supersonic plasmas and a planar solid obstacle gives rise to the formation of reverse shocks, propagating into the upstream as a result of material accumulation in the post-shocked region, near the obstacle. In the tungsten experiment, the structure of the reverse shock is characterised by a thin-shell configuration appearing as the brightest emitting profile in the stagnation region as

shown at $x = 0.7$ mm in Fig. 4.14(a). We refer to this sharp feature at $x = -1$ mm as "reverse shock" and its motion can be analysed from the sequence of self-emission images.

To estimate the speed of the reverse shock, the position of the shock can be tracked in each self-emission frame following a similar procedure used for the aluminium plasma. The shock feature appears as a dominant peak in the intensity profiles taken from horizontal line-outs at fixed z -position, as shown in Fig. 4.14(b). The distance between the obstacle ($x = 0$ mm) and the maximum peak can be considered as the width of the stagnated post-shocked region and serves as an estimate of the shock trajectory.

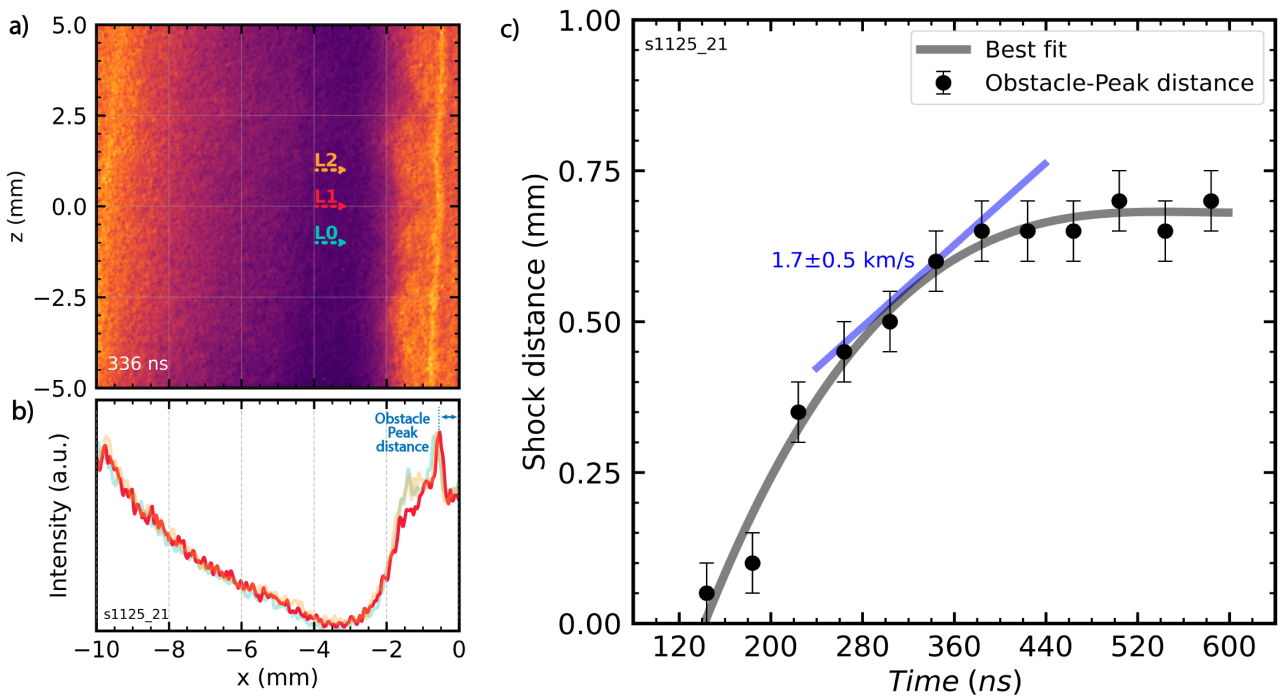


Figure 4.14: (a) An optical self-emission frame example captured at 336 ns after the current pulse. (b) Three intensity profiles from horizontal line-outs at various vertical positions, $z = 1$ mm, $z = 0$ mm and $z = -1$ mm. (c) Reconstructed trajectory of the shock expansion. The measurements of the shock distance from the obstacle are performed by taking as a reference point the peak maximum. The error bars are estimated by visually comparing the distance of the peak with the position of the shock.

The results of this analysis are presented in Fig. 4.14(c). The expansion of the shock from the obstacle surface follows a decelerating motion with an average velocity in the lab frame of $v_s = 1.7 \pm 0.5$ km/s and reaching $v_s \simeq 0$ km/s by $\simeq 400$ ns. This corresponds to an average deceleration in the lab frame of $a = 10^7$ km/s² under the assumption of constant deceleration. It is important to note that this hypothesis is an approximation, as the reverse shock motion is likely influenced by the

finite generator current pulse. The reconstructed motion of the reverse shock shows good agreement with a quadratic law during the expansion phase, however at a later time ($\simeq 450$ ns) after the current pulse, the shock distance plateaus, indicating that the assumption of constant deceleration is no longer valid. Another important parameter in the 2-D shock structure is the thickness. As more plasma flows

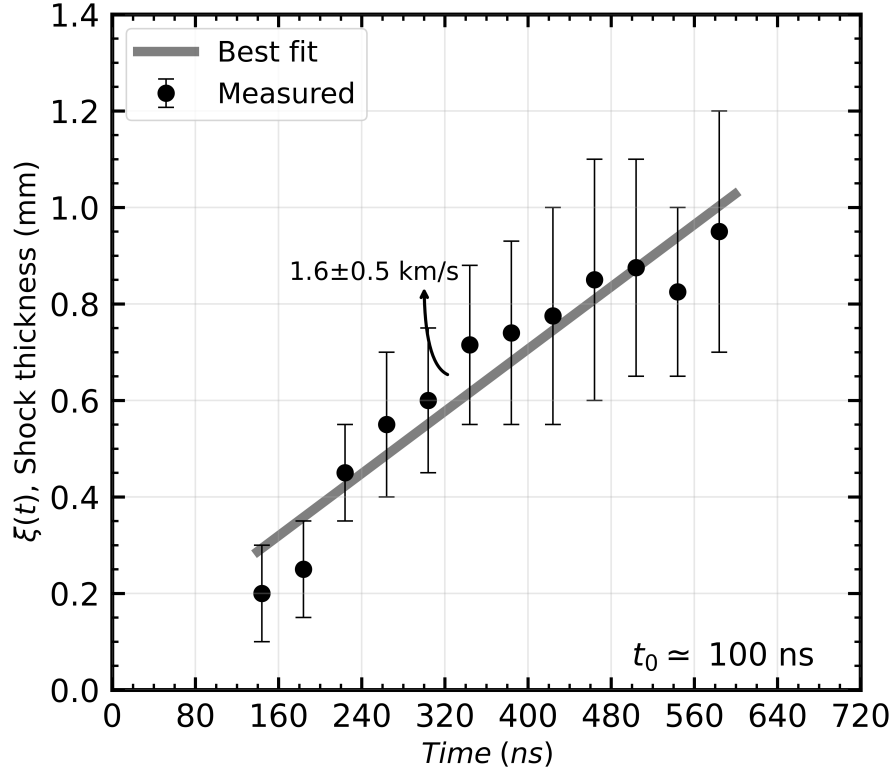


Figure 4.15: Time-resolved thin-shell thickness from optical self-emission images. Each data point represents the main peak width from horizontal intensity distribution across the shock feature. The measurements are performed by estimating the spatial range between 10% and 90% percentile of the cumulative intensity distribution of the main peak.

reach the reverse shock, a portion of the post-shocked material contributes to enlarging the thin-shell structure. The thickness of this thin shell can be estimated directly from horizontal intensity line-outs by measuring the width of the main peak feature, representing the reverse shock (Fig. 4.14). In each frame, the width of the layer is represented by the distance between the two local minima across the intensity peak, which is estimated as the spatial range between 10% and the 90% percentile of the cumulative intensity distribution of the main peak. The measured thickness profile over time is shown in Fig. 4.15. Starting from the initial frame captured at 144 ns after the start of the current, the thickness increases linearly over time, ranging from 200 μm to approximately 1 mm in the final frame. The rate of the layer thickness growth is estimated to be 1.6 ± 0.5 km/s, which is remarkably close to

the average reverse shock speed measured in Fig. 4.14(c). It is noteworthy to say that this thickness increase could also be an indication of the presence of perturbations in the layer, integrated along the diagnostic line-of-sight. The importance of these parameters will be discussed in appendix A.

The uncertainties in all the above measurements can be attributed primarily to depth of field errors. As we are imaging an extended thin object, the presence of less-than-perfect sharpness in the image can introduce errors in direct spatial measurements of feature size. A rough estimate of the error introduced by the depth of field can be achieved by comparing the measured thickness of the planar obstacle in a background image without a plasma ($\simeq 650 \mu\text{m}$) using the same procedure described earlier, with its known geometrical value ($\simeq 500 \mu\text{m}$). This reveals an error difference of around $150 \mu\text{m}$, corresponding to an uncertainty of approximately 23%. There are other sources of uncertainties when working with self-emission cameras. The optical camera used in this study consists of 12 CCD detectors, each with dimensions of 649×558 pixels (width $\times$ height), a physical size of 12×10 mm, and a pixel size of $18 \mu\text{m}$. The spatial resolution of the sensors is determined by dividing the pixel size by the magnification factor provided by the optical telescope, which ranges from approximately 0.3 to 0.6, resulting in a resolution of about 30 to $60 \mu\text{m}$. Another possible source of error could arise from the finite exposure time when imaging fast-moving objects. For moving features at speeds of approximately $\simeq 2$ km/s, an exposure time of a few nanoseconds (5 ns) would correspond to a displacement error of approximately $10 \mu\text{m}$. Both the resolution of the sensors and the exposure time error contribute much less to the overall uncertainty compared to the depth of field effect, thus their influence can be neglected.

Characterise the spatial non-uniformities of the reverse shock

Self-emission images are the only available time-resolved data that can be used to characterise the spatial non-uniformities along the thin shocked layer. Based on the measured shock speed and layer thickness, it is expected that the system may be susceptible to unstable waves with spatial scales resolvable by the optical self-emission diagnostic. To analyse the instability, two different approaches can be performed: contour displacement (not shown) and intensity profile. The spatial displacement method involved directly tracking the contour of the thin-shell layer, resulting in positions of constant

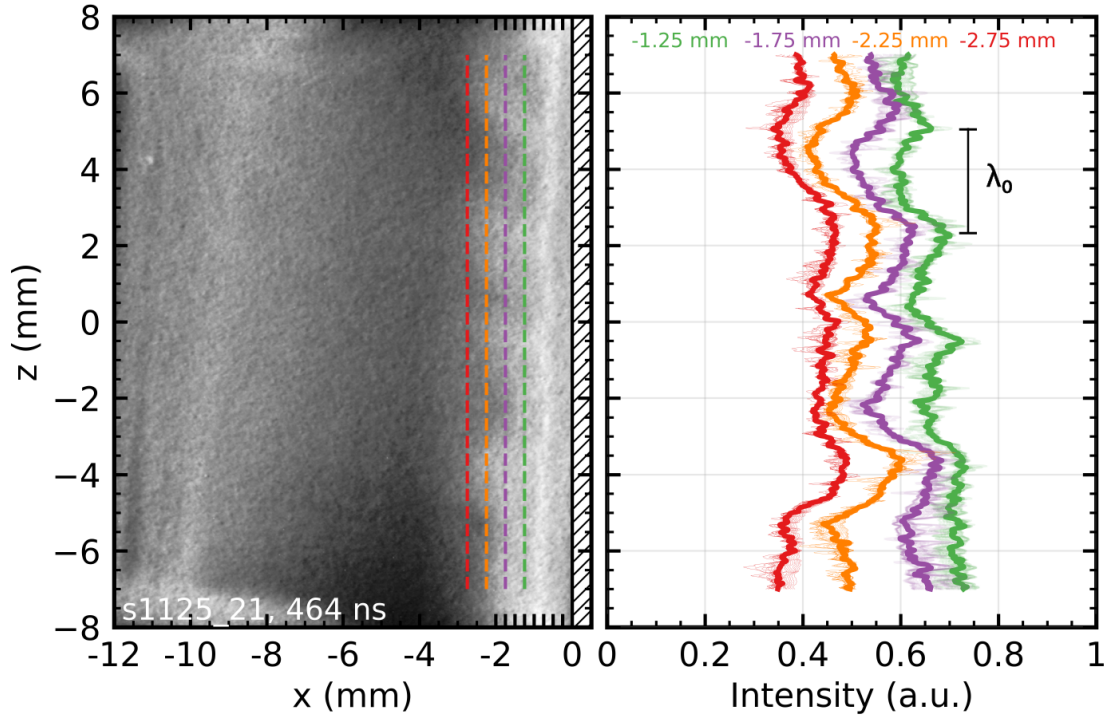


Figure 4.16: Example of vertical intensity profiles taken at 0.5 mm intervals in the range between $x = -1.25$ mm and $x = -2.75$ mm at 464 ns after the current pulse. The characteristic wavelength λ_0 reflects the longest modulation along the shocked layer.

intensity. The other approach relies on intensity profiles taken from vertical line-outs at a fixed distance from the shock boundary which can be used to identify the longest visible wavelength. While both methods yielded similar results, the intensity profile approach proved to be less noisy and more effective compared to the contour displacement method. The distribution of plasma materials surrounding the thin layer results in an intensity profile that reflects the corrugation of the layer, which can be seen in Fig. 4.16. Four different vertical line-outs taken at 0.5 mm intervals in the spatial x -axis demonstrate a similar intensity distribution in Fig. 4.16(Right). The bold lines indicate the computed median intensities from twenty different average thicknesses ranging from 0.05 mm to 2 mm (shown as shaded lines) which helps to reduce intensity fluctuations. The intensity profile exhibits a quasi-periodic nature with a characteristic wavelength estimated to be $\lambda_0 \approx 3 \pm 0.5$ mm. This corresponds to a spatial frequency of $k_{z,0} = 2\pi/\lambda_0 \approx 2 \pm 0.5 \text{ mm}^{-1}$. This characteristic wavelength does not change across frames where the non-uniformity is visible, $t > 300$ ns.

It is important to note that due to the limitations of the diagnostic, definitive statements cannot be made about shorter wavelengths. Therefore, we consider the observed mode to be monochromatic

representing the dominant mode along the thin-shell layer. Away from the shock front projection effects smooth out the observed modulation, leading to overestimated values of λ_0 , as expected.

4.3.2 Characterising the upstream flow in tungsten reverse shock

In the previous subsection, the analysis of the overall reverse shock dynamic has been performed from optical self-emission images. The results revealed the two-dimensional nature of the reverse shock represented by a bright narrow profile within the interaction region. This narrow shocked structure expands upstream towards the wire array, increasing in thickness and showing a well-defined modulation along its vertical length. Ahead of the reverse shock, visible perturbations are observed at late time ($t > 300$ ns) from the plasma self-emission with no clear sign of stand-off shocks like in the aluminium experiments (Fig. 4.12). Laser-based diagnostics such as interferometry, shadowgraphy, the novel imaging refractometer and optical Thomson scattering have been employed to further characterise the overall structure of tungsten reverse shock.

Shadowgraphy and imaging refractometer data

Overall shock structures in tungsten plasma flow interacting with a planar obstacle exhibit a significantly different behaviour compared to the aluminium case, despite having similar upstream mass density and flow velocity. Perturbations ahead of the reverse shock observed in the self-emission images suggest that these regions could be dominated by non-uniform plasma density. Under these conditions, the shadowgraphy and imaging refractometer techniques allow us to further characterise these observed structures.

The optical diagnostic setup described in § 4.1 is employed to image the side-on view (xz-plane) of the plasma interaction region as shown in Fig. 4.17. A background image was taken before the shot in the presence of a metallic pin placed at a known distance from the planar obstacle so that the physical scale of the image can be accurately determined. Here, the origin of the coordinate system is chosen such that the wires are located at approximately $x = -10$ mm, while the obstacle edge is at $x = 0$ mm. Fig. 4.17(b) shows the shadowgraphy image during the experiment, captured at 440 ns after the current

start. In the shot image, the wires are clearly visible at around 10 mm from the obstacle, possibly due to lower electron density than the aluminium wire array. It is important to note that the intensity of the

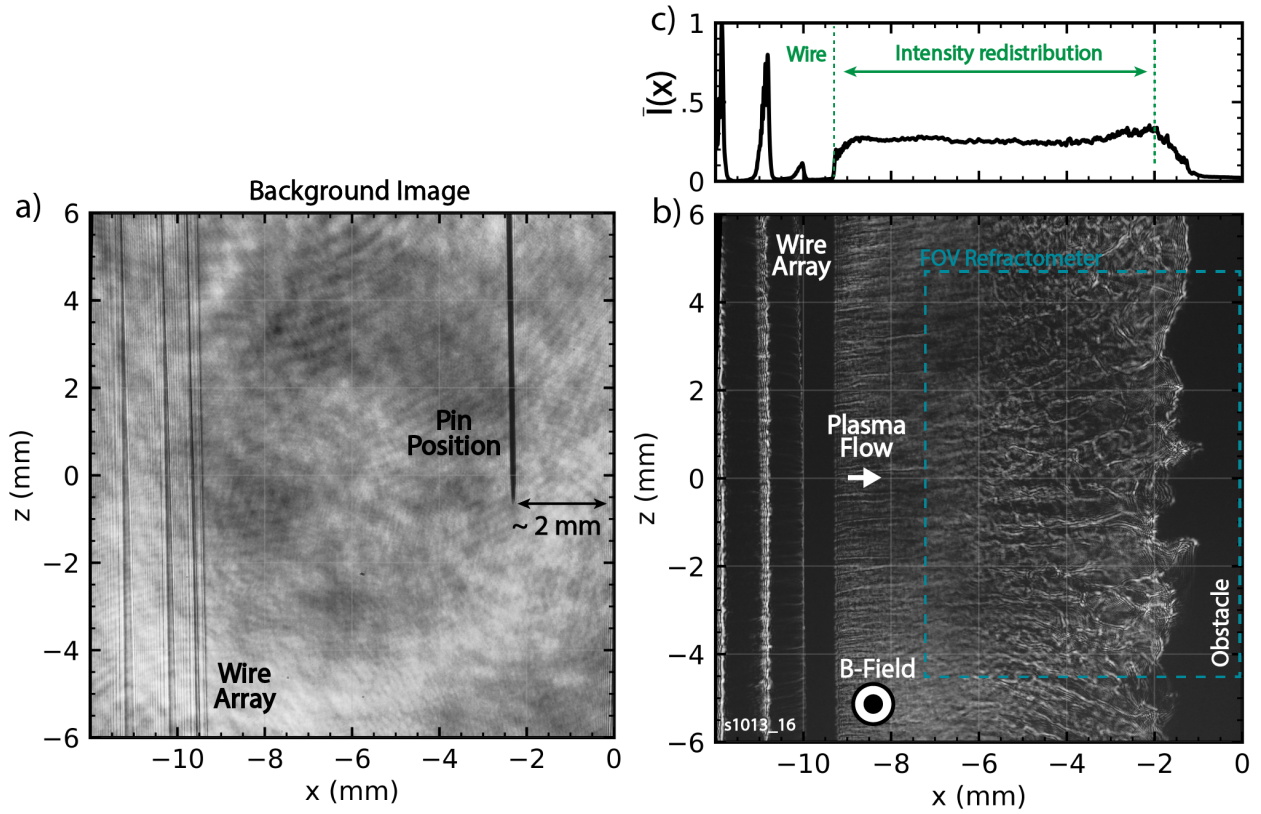


Figure 4.17: (a) Background image taken before the experiment. Only the front wires are visible at $x = -10$ mm with the obstacle edge placed at the origin of the x-axis. (b) 1053 nm side-on shadowgram of tungsten plasma flow colliding against a planar obstacle at 400 ns after the current start. The initial perturbation imposed by the wire ablation phase is visible in the upstream flow until $x = -6$ mm, while small-scale structures dominate the plasma flow up to the opaque region near the obstacle. (c) sum-up intensity profile, $\bar{I}(x)$. In the region between the wires ($x = -9$ mm) and the boundary of the stagnation layer, the laser intensity is almost constant. This indicates that the probing beam intensity is completely redistributed due to changes in the plasma refractive index without losses caused by strong laser probe deflection.

shadowgraphy image remains flat in the region between the wires, $x = -9$ mm, and the boundary of the stagnation layer, $x = -2$ mm [see Fig. 4.17(c)], indicating that in this region, the laser intensity is completely redistributed due to changes in the plasma refractive index without losses caused by strong laser probe deflection. Two distinct regions can be observed from the shadow image of the plasma flow between the wire array and the stagnated plasma at the obstacle. A relatively smooth plasma flow is observed near the wire array with pronounced small-intensity modulations oriented along the plasma flow direction, extending from $x = -9$ mm up to $x = -6$ mm. For $x > -6$ mm, the upstream intensity modulations become less pronounced, and small structures start to dominate the plasma flow. The

transition happens smoothly with no presence of abrupt intensity discontinuities, indicating the absence of a well-defined stand-off shock. In the region between $x = -6$ mm and $x = -2$ mm, the intensity distribution appears fragmented with modulations oriented both in z and x directions, showing a wide range of spatial scales. These small-scale perturbations do not seem to develop in the absence of the obstacle. Near the obstacle ($x > -2$ mm), the interaction region appears opaque with visible ripples at the region interface. The imaging refractometer technique, described in Chapter 3, is employed along

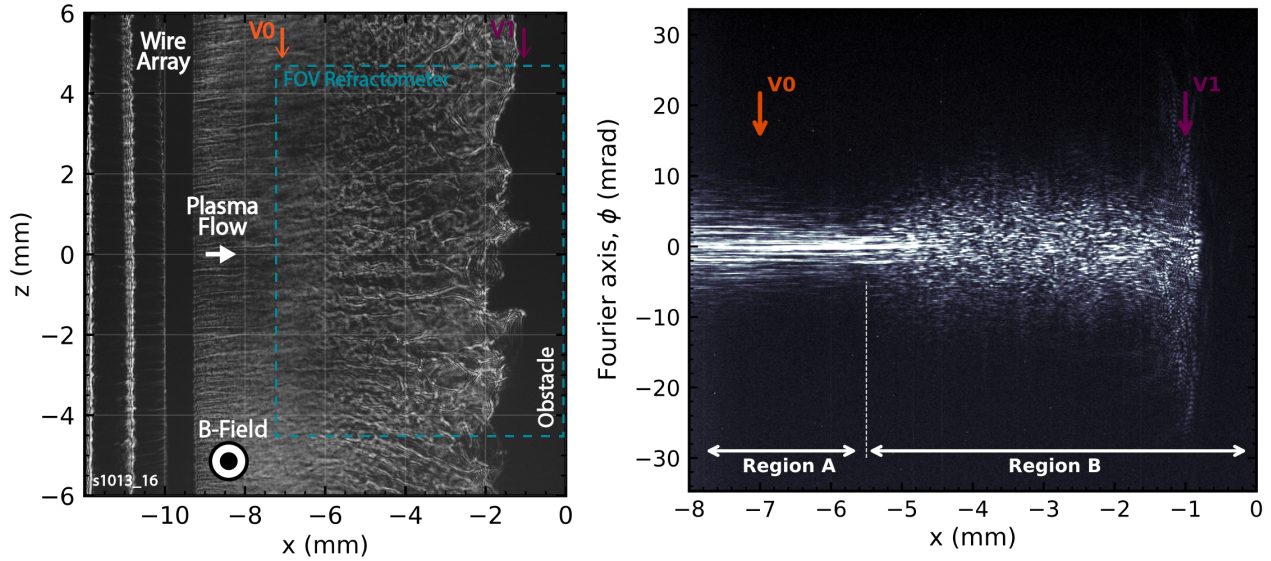


Figure 4.18: (Left) shadowgram presented in Fig. 4.17. (Right) refractogram of the FOV region highlighted in the shadowgraphy image both taken at the same time. Notice that the entire precursor region ahead of the obstacle is characterised by small-scale intensity structures, indicating that the laser probe propagates through an inhomogeneous media. Two vertical line-outs labelled V0 and V1 are taken at $x = -7$ mm and $x = -1$ mm resulting in the angular intensity distributions of Fig. 4.19.

with shadowgraphy to gain more quantitative insights into the perturbations developed in the plasma, ahead of the stagnation layer. The refractogram is shown in Fig. 4.18 and the field of view (FOV) of the refractometer corresponds to the region highlighted in the shadow image above. Two distinct regions, similar to those seen in the shadowgraphy data, can be identified as Region A, extending from the wire array at $x = -10$ mm to $x = -5.5$ mm, and Region B, spanning from $x = -5.5$ mm to the obstacle surface at $x = 0$ mm. For both regions, there is a significant broadening of the angular distribution in comparison to the non-perturbed probing beam.

Similar to the aluminium case, in Region A there is a reduction in the width of the deflection angle distribution as the plasma flow expands outward from the array. The angular distribution is modulated

by the interference effects arising from the wire-array ablation perturbations, similar to those seen for aluminium, but at slightly different periodicity. In Region B, the distribution of deflection angles widens, reaching approximately 7 mrad over a broad region. In close proximity to the obstacle, the distribution expands further and appears stronger, characterised by small-scale fluctuations in both the x and ϕ directions. As already seen from the shadowgraphy image, the transition between the two regions, A and B, is smooth with no abrupt intensity cut-off, in strong contrast to the aluminium refractogram (Fig. 4.8). Characteristic angular deflection distributions for Region A and Region B

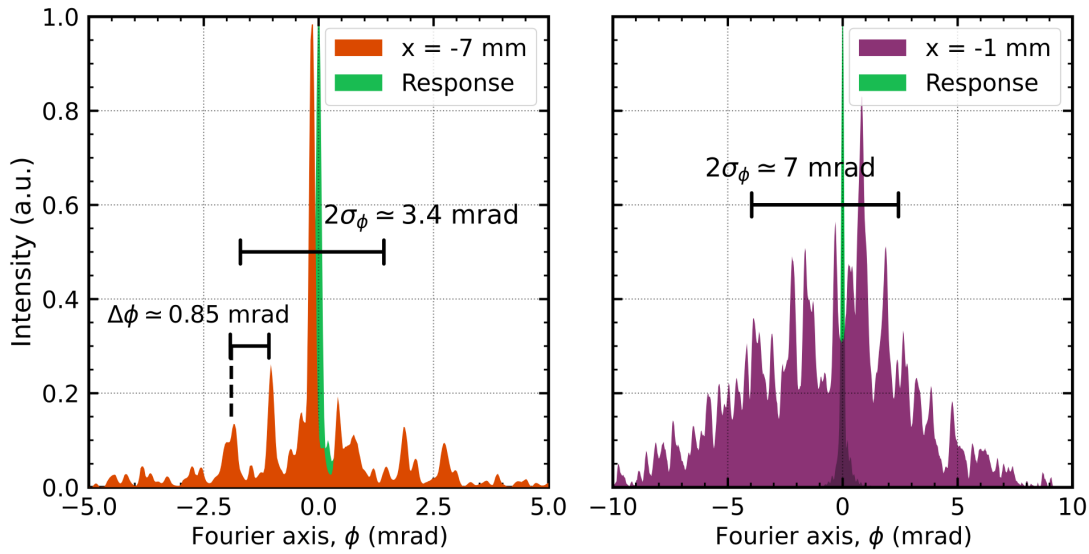


Figure 4.19: (a), (b) Example of angular deflection intensity profiles at $x = -7$ mm and $x = -1$ mm from the obstacle, respectively. After the transition from region A to region B, the density gradients increase together with the angular broadening, σ_ϕ . Peak displacements resulting from interference effects are equally spaced with $\Delta\phi \simeq 0.85$ mrad, visible in the upstream flow.

are shown in Fig 4.19. The laser intensity profile has a width of $2\sigma_\phi \approx 3.4$ mrad and also has evenly spaced peaks separated by $\Delta\phi \approx 0.85$ mrad at $x = -7$ mm. These periodic peaks are similar to those seen in aluminium but with a higher frequency corresponding to a shorter wavelength (≈ 1.2 mm) of density perturbations in the flow than what was previously observed in aluminium plasmas. In Region B, for the angular profile at $x = -3$ mm, the intensity profile widens significantly, with $2\sigma_\phi \approx 7$ mrad. The angular profile is highly modulated but cannot be described as having a single periodicity. The spatial evaluation of the angular distribution, measured as a standard deviation σ_ϕ of the distribution is shown in Fig. 4.20. It is seen that the width (standard deviation) decreases to a minimum value of $\sigma_\phi \approx 1$ mrad at $x = -5.5$ mm, at the edge of Region A. The transition between regions A and B is marked by a sharp increase in the width of the distribution of angular deflections from $\sigma_\phi \approx 1$

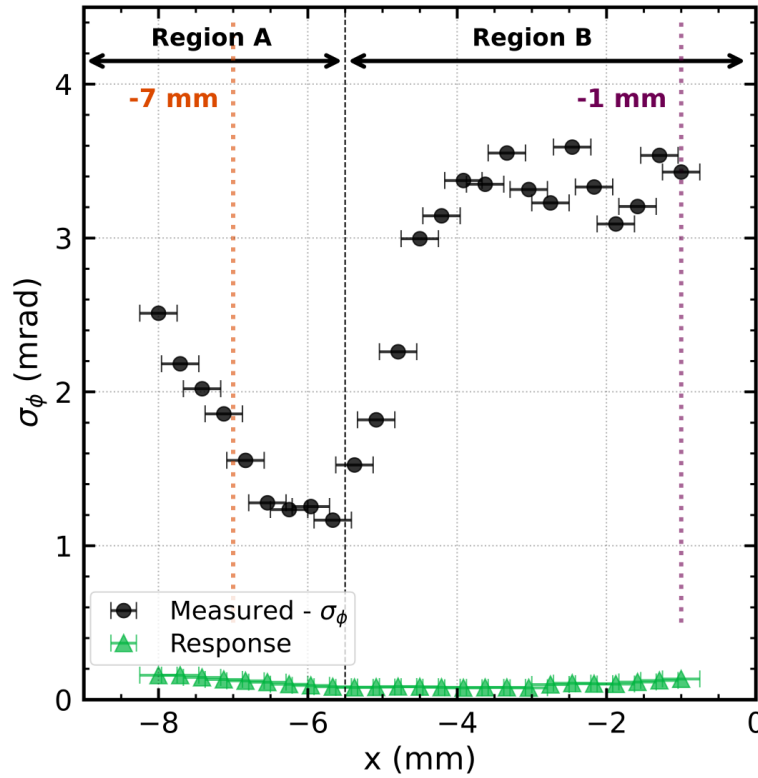


Figure 4.20: The standard deviation of the refractogram signal, σ_ϕ . The error bars represent the spatial averaging of the vertical line-outs, $\Delta x = 0.4$ mm.

mrad to $\sigma_\phi \approx 3.5$ mrad, between $x = -5.5$ mm and $x = -4$ mm. For $x > -4$ mm, the angular deflection signal remains relatively constant at approximately the same values of 3.5 mrad. It is important to note that at all positions in the refractometry image, the angular deflection is significantly greater than the diagnostic response (indicated by the green line), with $\sigma_{\phi,\text{signal}}/\sigma_{\phi,\text{response}} \approx 200$.

Electron Density Measurements

The laser interferometry technique is widely used in high energy density plasma to measure the line-integrated electron density $\int n_e dl$, as already introduced in chapter 2. Exploding wire arrays with tungsten wires have not been extensively studied in the past; therefore, both Thomson scattering and electron density measurements have been performed in tungsten experiments with a planar obstacle along with a free-flow experiment where the obstacle was removed from the system. In both experiments presented in this section, a 532 nm laser imaging interferometry is used to image the plasma from the end-on view (see setup schematic 4.2). Specifically, the laser beam is brought into the plasma along the z -direction, so that the imaging takes place in the xy -plane of the plasma interaction.

Accurate measurements of electron density are essential for characterising the plasma state in the upstream region. Fig. 4.21 shows the end-on raw interferograms obtained using a 532 nm laser probe.

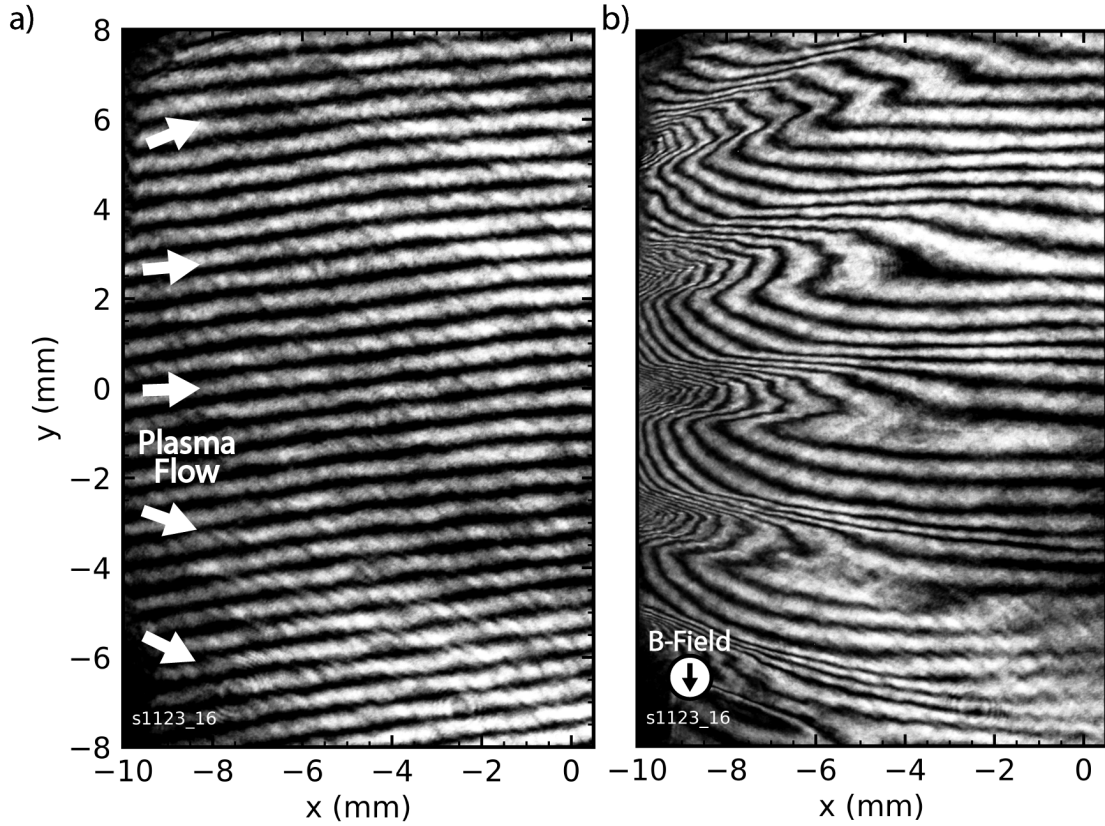


Figure 4.21: (a) Background end-on interferogram of the null-shot (s1123_16). The edge of the wire array is located at $x = -10$ mm. (b) End-on shot interferogram taken at 384 ns after current pulse in the presence of plasma.

The wire array is located outside the field of view, on the left-hand side of the image, so the plasma outflow expands radially from left to right, starting at $x = -10$ mm. The background interferogram was taken before the experiment in the absence of plasma and is shown in Fig. 4.21(a). The fringe spacing sets the resolution of the diagnostic while the orientation of the fringes is chosen with a small initial angle so that the measurement is sensitive to both axes. In the shot interferogram reported in Fig. 4.21(b), distorted fringes are observed due to the presence of plasma during the experiment. Fringes can be traced throughout the diagnostic field of view with a few regions of lost fringes observed near the wires (at $x = -10$ mm), caused by the probing beam being deflected by large density gradients or absorbed by inverse bremsstrahlung. The result of the interferogram analysis described in chapter 2 is presented in Fig. 4.22(a) as the line-integrated electron density map, $n_e L$. The corresponding volumetric electron density can be obtained by assuming the integration length equal to the height of

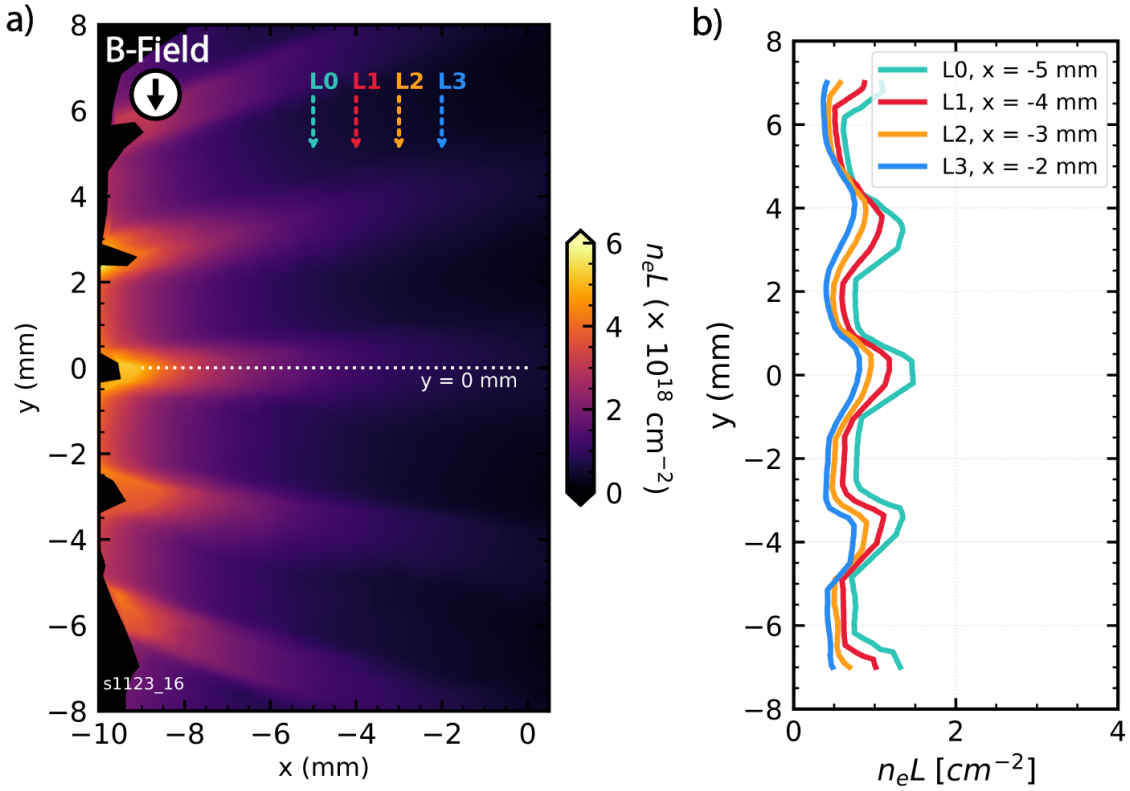


Figure 4.22: (a) Electron density map reconstructed from the raw interferograms presented Fig. 4.21. The shot interferogram was taken at 384 ns after the current pulse during the free-flow experiment (s1123_16) in the absence of any obstacles. (b) Electron density profiles, $n_e L$, from three vertical line-outs at $x = -5 \text{ mm}$, $x = -4 \text{ mm}$, $x = -3 \text{ mm}$.

the wire array, $L = 22 \text{ mm}$. It is seen from Fig. 4.22(a) that the highest electron density is recorded near the wire array at $x = -10 \text{ mm}$ and density decreases away from the array. Azimuthal density modulations are clearly visible and regions of higher density are separated by regions of lower density. Areas of high density are approximately conical in shape and are aligned with gaps between the wires. These are the result of the formation of oblique conical shocks in which plasma outflows from adjacent wires collide, similar to those observed in the ablation flows in standard wire arrays [69, 86, 87, 89, 103].

The measured electron density is $n_e \simeq 2 \times 10^{18} \text{ cm}^{-3}$ near the wire array and decreases along the plasma flow to a minimum density of $n_e \simeq 0.5 \times 10^{18} \text{ cm}^{-3}$ at 10 mm from the array. Examples of density profiles taken along vertical line-outs at three positions, $x = -5 \text{ mm}$, $x = -4 \text{ mm}$, $x = -3 \text{ mm}$, are shown in Fig. 4.22(b). The azimuthal modulation of density is larger near the plasma source and flattens at larger distances. At $x = -2 \text{ mm}$ the modulation has a density ratio between minimum and

maximum of $\simeq 1.5$. At 10 mm from the array, the electron density distribution is expected to be almost uniform in the xy -plane.

Similar electron density measurements have been performed in experiments with the presence of a conductive planar obstacle. The obstacle was placed 10 mm from the wire array, orthogonal to the flow direction. Raw interferograms were taken at a comparable time of the null shot, 348 ns after the current start. Fig. 4.23(a) shows the raw end-on interferogram taken during the experiment. The wire array is located on the left-hand side of the image field of view, while the origin of the horizontal axis is at the obstacle position. It is seen that interference fringes are in the plasma flow for up to $x = -2$ mm. Closer to the obstacle, for $x > -2$ mm, fringes are lost in the interaction layer, most probably due to large density gradients.

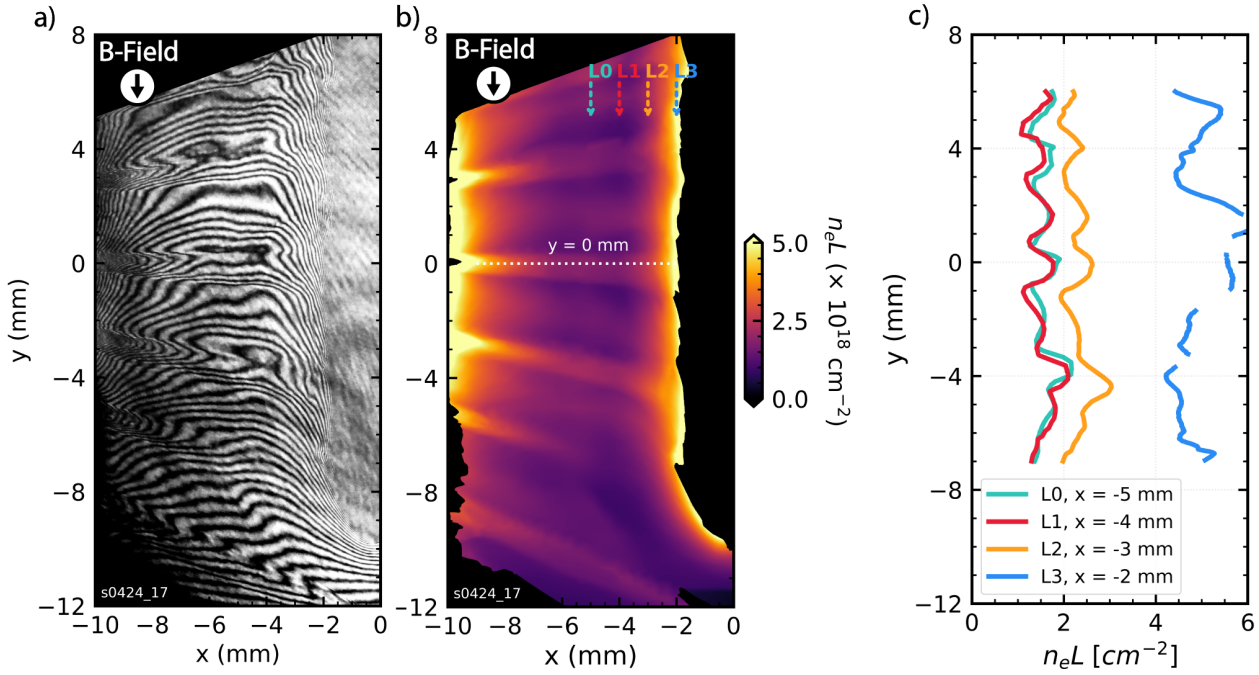


Figure 4.23: (a) End-on shot interferogram (s0424_17) taken at 348 ns after the current pulse. The planar obstacle is not visible and is located at $x = 0$ mm, outside of the data field of view. (b) Line-integrated electron density map. The magnetic field lines are transverse to the flow direction. (c) Density profiles taken from three different line-outs at $x = -5$ mm, $x = -4$ mm and $x = -3$ mm.

A region of zero fringe shift is available at $y = -10$ mm near the edge of the obstacle, allowing one to compute the absolute value of the line-integrated electron density. Fig. 4.23(b) shows the unfolded line-integrated electron density maps $n_e L$. The azimuthal density modulation is still observed in the experiment with the planar obstacle. In the upstream region, $x < -6$ mm, electron density

measurements are consistent with the free-flow shot, reaching a minimum density of $n_e \simeq 1 \times 10^{18} \text{ cm}^{-3}$ as the plasma expands towards the obstacle. The electron density starts to increase (in comparison to the free-flow case, see Fig. 4.22) at $x = -5 \text{ mm}$. This can be interpreted as due to the pile-up of the magnetic field advected by the flow ahead of the obstacle, similar to aluminium experiments. The maximum density of $n_e \simeq 3 \times 10^{18} \text{ cm}^{-3}$ is reached at the boundary of the stagnation layer, $x = -2 \text{ mm}$. For positions $x > -2 \text{ mm}$, the density fringes are lost due to strong density gradients.

Examples of density profiles from vertical line-outs are taken at the same positions as the null shot (Fig. 4.23) and are shown in Fig. 4.23(c). The data reveals that an obvious upstream modulation still persists at a large distance from the obstacle similar to those observed in the free-flow experiment, but with a much smaller amplitude. This azimuthal density modulation reduces closer to the obstacle, starting from $x = -2.5 \text{ mm}$ up to the boundary of the stagnation layer, $x = -2 \text{ mm}$, which appears uniform along the y -axis, parallel to the magnetic field lines. Although we cannot exclude the hypothesis that this modulation contributes to seeding the perturbations observed in the side-on self-emission images, it is unlikely that could be significant in the overall dynamic of the system. Interferometry

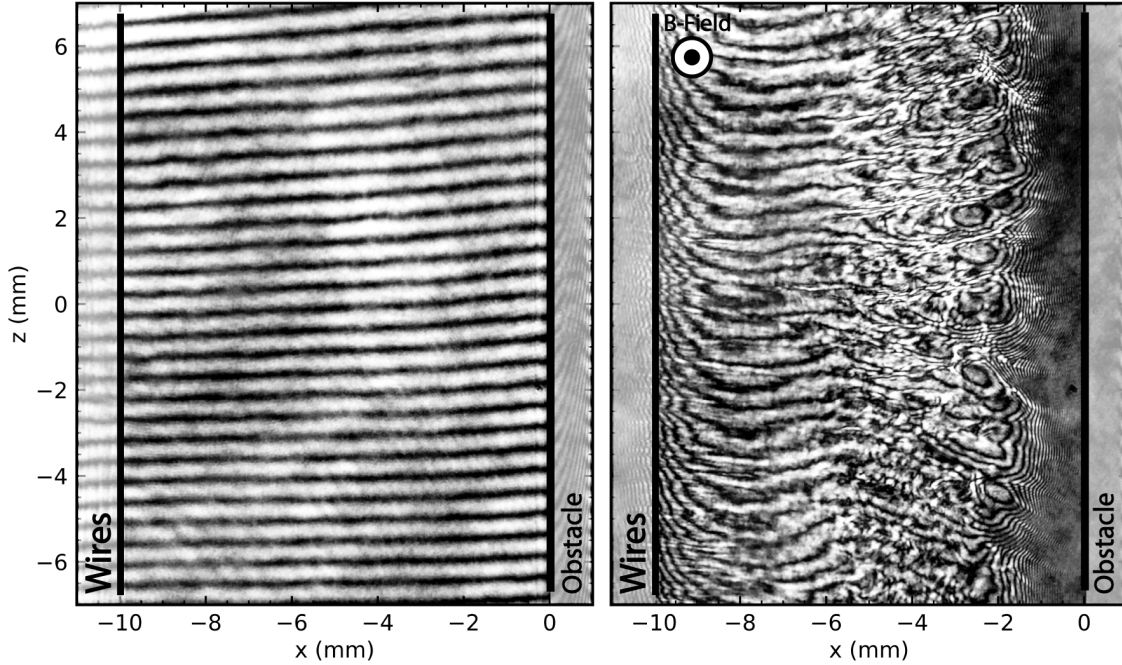


Figure 4.24: Side-on background (left) and shot (right) 1053 nm interferograms of tungsten flow-obstacle interaction. The front wires are visible at $x = -10 \text{ mm}$ with the obstacle edge placed at the origin of the x -axis. The shot interferogram was taken at 400 ns after the current start. Visible small-scale structures dominate the entire region ahead of the obstacle, especially for $x \geq -6 \text{ mm}$, not allowing to obtain the line-integrated electron density map.

measurements from the side-on view have been also performed using a 1053 nm laser probe directed along the y-direction, as shown in Fig. 4.24. The front wires are positioned at approximately $x = -10$ mm with the obstacle located at the origin of the horizontal axis. The raw interferogram is shown in Fig. 4.24(b) taken at 400 ns after the current pulse. Traceable fringe displacements are only observed near the wire array, while strong density perturbations become dominant starting from $x = -6$ mm. For $x \geq -6$ mm, fringes cannot be traced for the entire region ahead of the obstacle, not allowing to obtain more information than what was observed from the shadowgraphy data, Fig. 4.17(b). As already discussed previously, two regions of unperturbed and perturbed plasma flows can be observed from side-on images with a boundary at $x = -6$ mm. This is the same location where electron density deviates from the free-flow condition in the end-on density maps. In the perturbed region, $x \geq 6$ mm, the side-on data reveals strong non-uniformities on spatial scales smaller than those observed from the end-on view, which can indicate the presence of only k_z perturbations.

Optical Thomson Scattering Measurements

Thomson scattering diagnostic allows us to obtain localised measurements of plasma flow velocity, ion temperature and electron temperature by measuring the Doppler shift and fitting the spectral shape of the Thomson scattering signal, operating in the collective regime ($\alpha \simeq 1.6 \div 2.3$) and focusing on the observation of the ion acoustic feature (see chapter 2).

The Thomson scattering experimental setup designed to characterise the in-plane properties of both the upstream flow and the stagnation region is illustrated in Fig. 4.25. The same Thomson scattering configuration is applied for tungsten experiments without the planar obstacle. A focused laser beam propagates along the line $y = z = 0$ mm in the direction indicated by the wavevector $\mathbf{k}_{\text{in}}$. The scattered light is collected at 90° angle relative to the laser beam in the xy-plane from 14 equally spaced positions using a lens which images the probing laser beam onto a linear array of optical fibres. The overall magnification of the imaging system is 0.5, resulting in collection volumes of $\simeq 150 \mu\text{m}$ in diameter spaced by 0.3 mm. The scattering light from fibre-optic bundles is then recorded using an imaging spectrometer with a spectral resolution of $\sigma_\lambda = 0.25 \text{ \AA}$ and a temporal resolution set by the gate time ($\simeq 4$ ns). Usually, two fibre bundles placed in opposite directions are required to resolve the

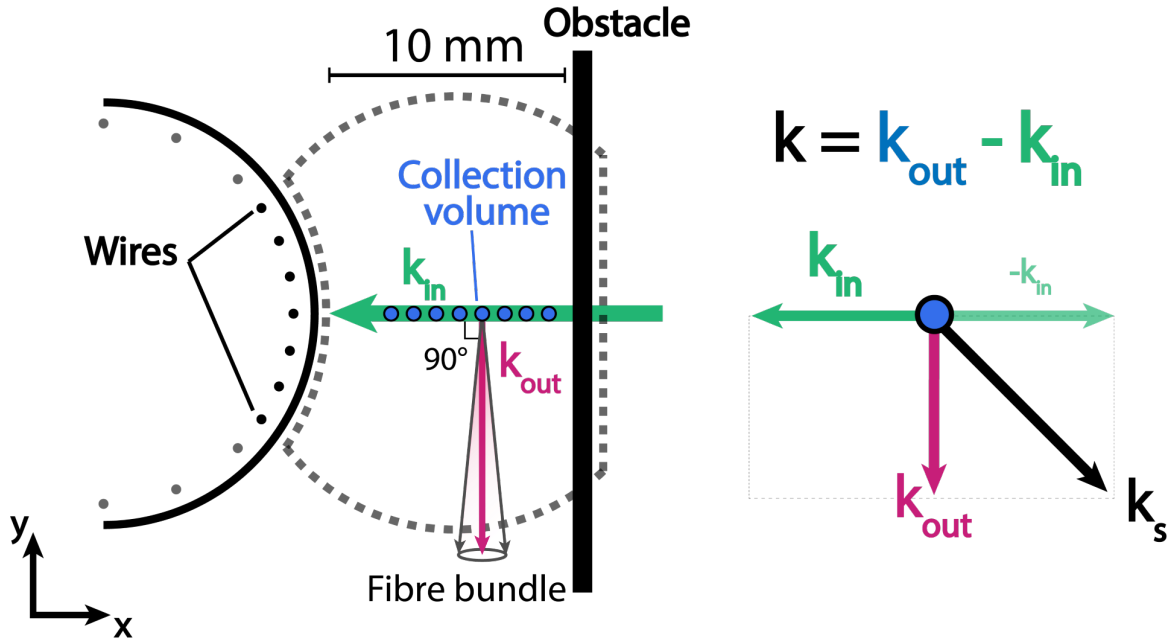


Figure 4.25: (Left) Schematic illustration of the Thomson scattering setup (not to scale). The 532 nm laser beam is focused and propagates through the upstream flow from the obstacle towards the array. To accommodate the laser beam both the obstacle and return post had pass-through holes. The laser beam is directed along the wavevector $\mathbf{k}_{in}$ while the scattered light is collected by a single fibre bundle positioned in the xy -plane at a 90° angle relative to the laser beam. Eight out of the 14 scattering volumes along the path of the laser are shown. The separation between each scattering volume is $\simeq 300\mu\text{m}$. (Right) Vector diagram of the Thomson scattering geometry. The resulting vector, $\mathbf{k}_s$, allows measurement of the plasma flow velocity vector in the xy -plane.

full in-plane velocity, however, preliminary experiments demonstrated the negligible contribution of the $\mathbf{V}_y$ component along the magnetic field, so that $\mathbf{V} \simeq \mathbf{V}_x$, and measurements were performed using only one single fibre bundle.

The described Thomson scattering geometry defines the vector diagram illustrated in Fig. 4.25(Right). The collection direction is represented by the vector $\mathbf{k}_{out}$ which is orthogonal to the laser beam $\mathbf{k}_{in}$. The resultant scattering vector is therefore oriented at 45° angle relative to the incident laser beam, measuring plasma flow velocities along the wavevector $\mathbf{k}_s = \mathbf{k}_{out} - \mathbf{k}_{in}$, from the Doppler shift of the Thomson scattering spectra. The fitting of the spectra also provides the thermal properties of the plasma, such as ion temperature and the product ZT_e . The present setup provides information on the spatial distribution of plasma parameters ($\mathbf{V}$, T_i , ZT_e) for a fixed time determined by the gating time of the camera.

Flow velocity and spectral broadening measurements in the plasma without obstacle

The Thomson scattering measurements of the tungsten free-flow were performed in the same experiment and at almost identical times ($t_{TS} = 365$ ns, $t_{interf} = 384$ ns) as the interferometry measurements presented in Fig. 4.22. The position of each Thomson scattering volume with respect to the electron density map was determined using the technique previously described in chapter 2 and are overlaid on the line-integrated density map as shown in Fig. 4.26. These measurements span over an $\simeq 8$ mm spatial range along the x-direction, approximately at the centre of the wire array, $z = y = 0$ mm.

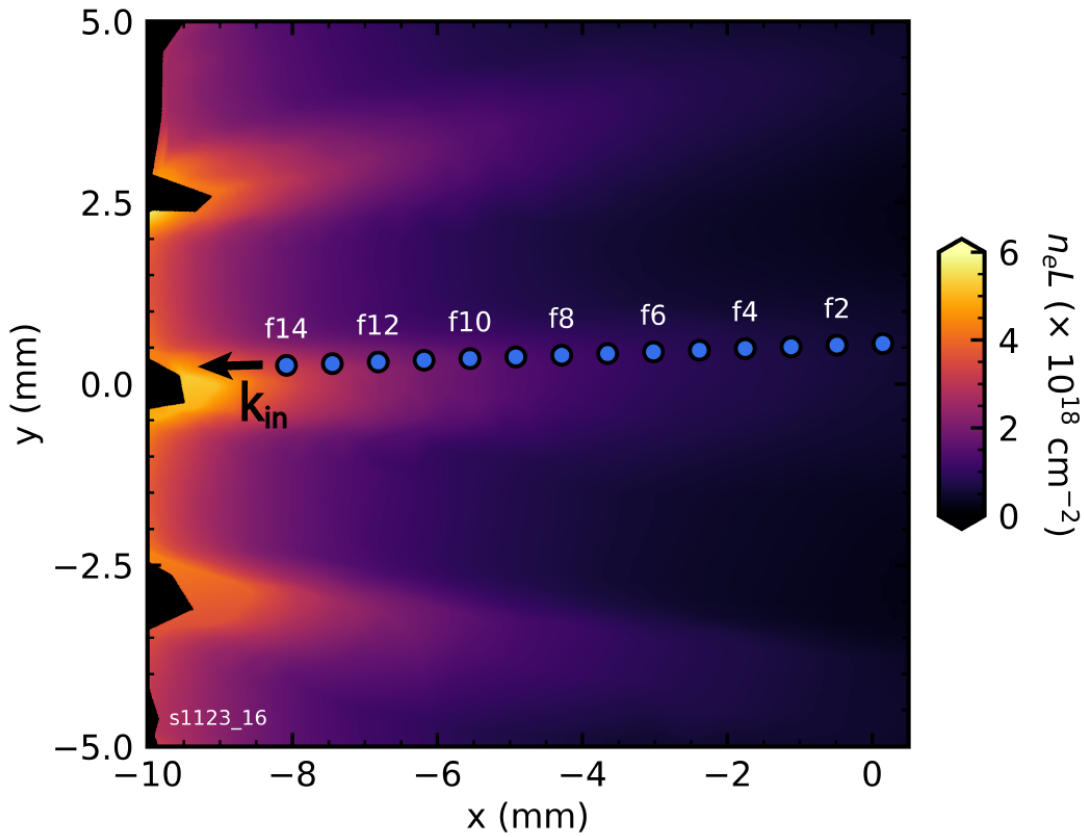


Figure 4.26: End-on electron density map from the tungsten null-shot (s01123_16) taken at 384 ns after the current start. The location of each scattering volume is determined by interpolating the positions of the scattered light of a metallic object along the laser light in pre-shot images.

The collection volumes are numerically labelled from 1 to 14 as shown in Fig. 4.26 (also in Fig. 4.27). The fitting process for each collected spectrum is constrained by the value of electron density at each corresponding location obtained from the interferometry measurements.

The raw spectrogram collected from this experiment is presented in Fig. 4.27 together with an

example of the Thomson scattering spectrum. It can be seen from the spectrogram in Fig. 4.27(a), that the signals are concentrated within a $2 \text{ \AA}$ range near the 532 nm laser wavelength, as expected for the ion-acoustic feature of collective Thomson scattering. Overall, spectral signals are detected from all

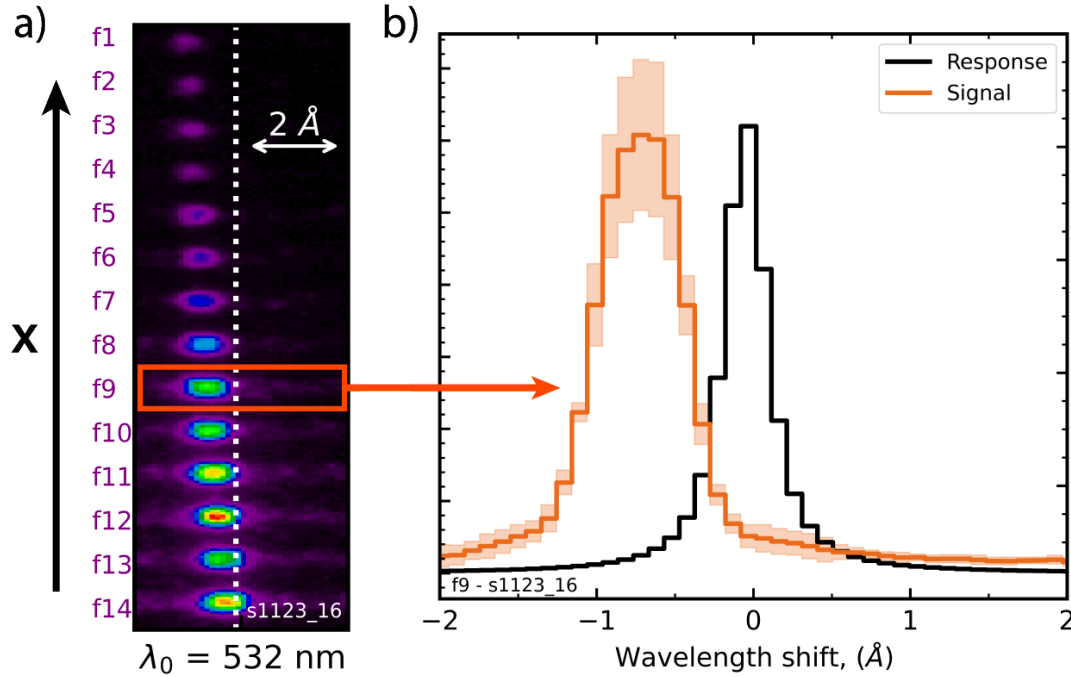


Figure 4.27: (a) Thomson scattering spectrogram (s1123_16) recorded at 365 ns after the current pulse. The laser wavelength is indicated by the dotted white line. Weak signals are recorded in fibres 1 to 7. Doppler shift of the spectra reduces from fibre 1 to 14. (b) Example of Thomson scattering spectrum from fibre 9 ($x = -5.3 \text{ mm}$). The shaded region represents the standard deviation of the signal while the solid line is the average over one fibre in the vertical direction of the spectrogram. The black line represents the response of the Thomson scattering diagnostic in the absence of plasma.

the available collection volumes with no evidence of the characteristic double-peak of the ion acoustic waves. The recorded signals from fibre 1 to 7, positioned further from the wire array, are weaker than the remaining scattering volumes, fibre 7 ÷ 14 which could be due to the decrease in plasma density. The spectral intensity of plasma self-emission, extending over a wide range of wavelengths outside the ion acoustic spectra, is visible in fibres 9 ÷ 14.

For fibres 1 ÷ 7, corresponding to $x > -4 \text{ mm}$, spectra exhibit significant Doppler shift relative to the probing laser wavelength. The blue shifts of the Thomson scattering signals are consistent with the direction of the plasma flow radially outwards. Fibres from 8 to 1 show much higher signal strength (consistent with higher plasma density), while the blue shift is smaller, consistent with plasma being accelerated in this region.

In addition to the Doppler shifts, there is some broadening of the Thomson scattering spectra, slightly exceeding the resolution of the spectrometer, especially closer to the wire array. An example of Thomson scattering spectrum from fibre 9 is shown in Fig. 4.27(b). The Thomson scattering signal is represented by the solid line obtained by summing up the intensity over one fibre in the vertical direction of the spectrogram to achieve a higher signal-to-noise ratio against self-emission. The shaded region represents the standard deviation of the signal. The spectral response function of the diagnostic ($\sigma_\lambda \simeq 0.3 \text{ \AA}$) is shown in solid black centred at the 532 nm laser wavelength and determined by measuring the spectrum of the scattered light of the probe beam from a thin metallic object, as described in chapter 2. It is important to note that the width of the spectral response function is only a factor of 2 or so smaller than the width of the Thomson scattering spectra for tungsten which poses a challenge for measurements of the plasma temperature. The results from the Doppler shift

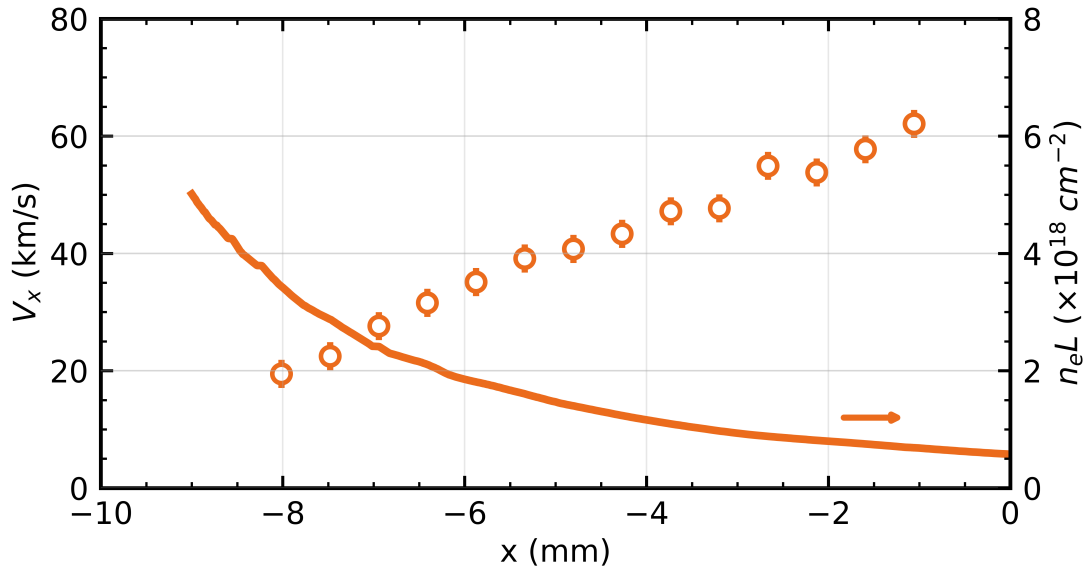


Figure 4.28: (Circle markers) Velocity component V_x from the measured Doppler shift of the Thomson scattering spectra. The error bars are small and fall within the marker size of the data. (Solid line) Line integrated electron density taken from a horizontal line-out at $y = 0$. The spatial range covered by the scattering volumes is 7 mm.

analysis of the Thomson scattering spectrogram recorded in the free-flow experiment are presented in Fig. 4.28. The specially resolved velocity component V_x is plotted together with the line-integrated electron density profile along the horizontal line-out $y = 0$ mm in Fig. 4.22. The recorded velocity starts at 20 km/s at $x = -8$ mm and increases linearly up to the maximum velocity of 60 km/s at $x = -1$ mm. The uncertainties of Doppler shift measurements are very small, $\simeq 1 - 2$ km/s, producing

errors that are well within the marker size. As already mentioned, the component V_y is negligible and therefore it is not measured in these current experiments. The measured velocity profile follows the inverse of the electron density profile $n_e L$, as expected from mass conservation.

Although reliable measurements of the thermal properties of tungsten plasma cannot be obtained from the current Thomson scattering data, the spectral broadening observed in scattering spectra near the wire array (fibre 7÷14) can be used to estimate an upper limit on the expected temperatures. A simple way to estimate the broadening of the Thomson scattering spectra is to calculate the intensity-weighted standard deviation of the Thomson scattering signal, knowing that $\Delta\lambda \simeq 2\sigma_\lambda$, where σ_λ is calculated as follows:

$$\sigma_\lambda = \frac{\sqrt{\sum_i \lambda_i^2 I(\lambda_i)}}{\sum_i I(\lambda_i)} \quad (4.1)$$

where $I(\lambda_i)$ is the pixel intensity of the spectrum. This quantity represents a statistical description of the spectral width without making assumptions about the intensity distribution. It is worth noting that this approach is sensitive to the background noise introduced by the plasma self-emission. For this reason, the background self-emission (outside of the ion-acoustic spectral feature) is averaged out and subtracted from each raw Thomson scattering spectrum before computing σ_λ . An example of this procedure is presented in Fig. 4.29. The result of the spectral width analysis performed for each

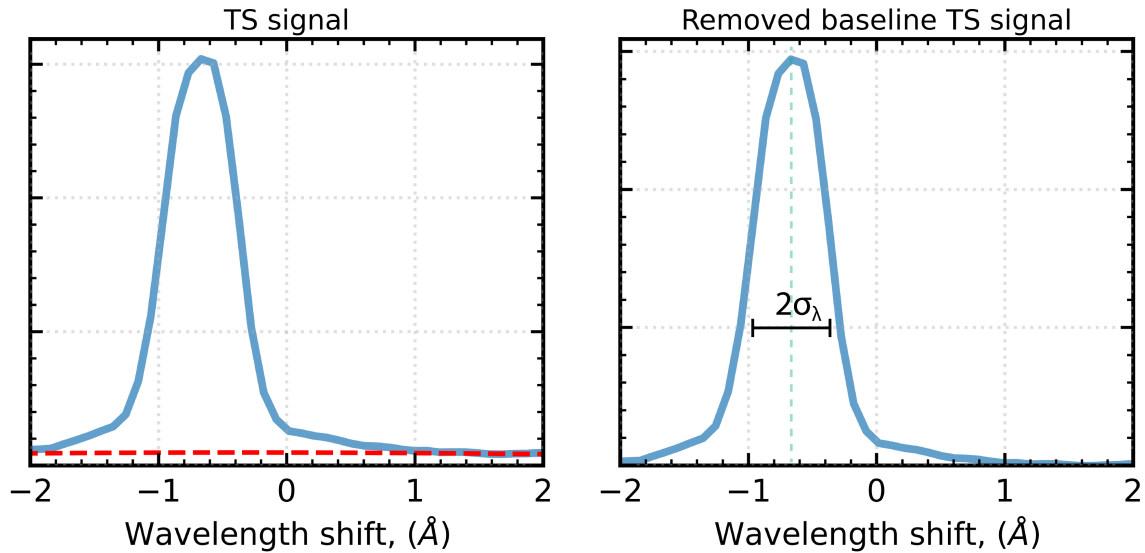


Figure 4.29: Example of raw Thomson scattering spectrum after the subtraction of the background self-emission. The spectral width is estimated as $2\sigma_\lambda$.

Thomson scattering spectrum is summarised in Fig. 4.30(Top), together with the value of σ_λ measured

from the diagnostic response spectra. The width of the diagnostic response function is almost constant across each scattering volume with an average value $\sigma_\lambda \simeq 0.3 \text{ \AA}$. For spectra collected near the wire array, $x < -4 \text{ mm}$, the spectral width recorded in the free flow experiment is only 40% wider than the diagnostic response width, $\sigma_\lambda \simeq 0.45 \text{ \AA}$. Despite being small, the difference between the signal and the diagnostic resolution is way larger than the potential broadening effects caused by noise, therefore an upper limit on temperatures can be estimated. In contrast, scattering volumes further away from the plasma source, $x > -4 \text{ mm}$, exhibit smaller spectral width, $\sigma_\lambda \simeq 0.35 \text{ \AA}$, in comparison to the diagnostic response function, hence no information can be inferred from these data points. The error bars of the measurements are estimated based on the uncertainties of the Thomson scattering spectra, as shown in Fig. 4.27.

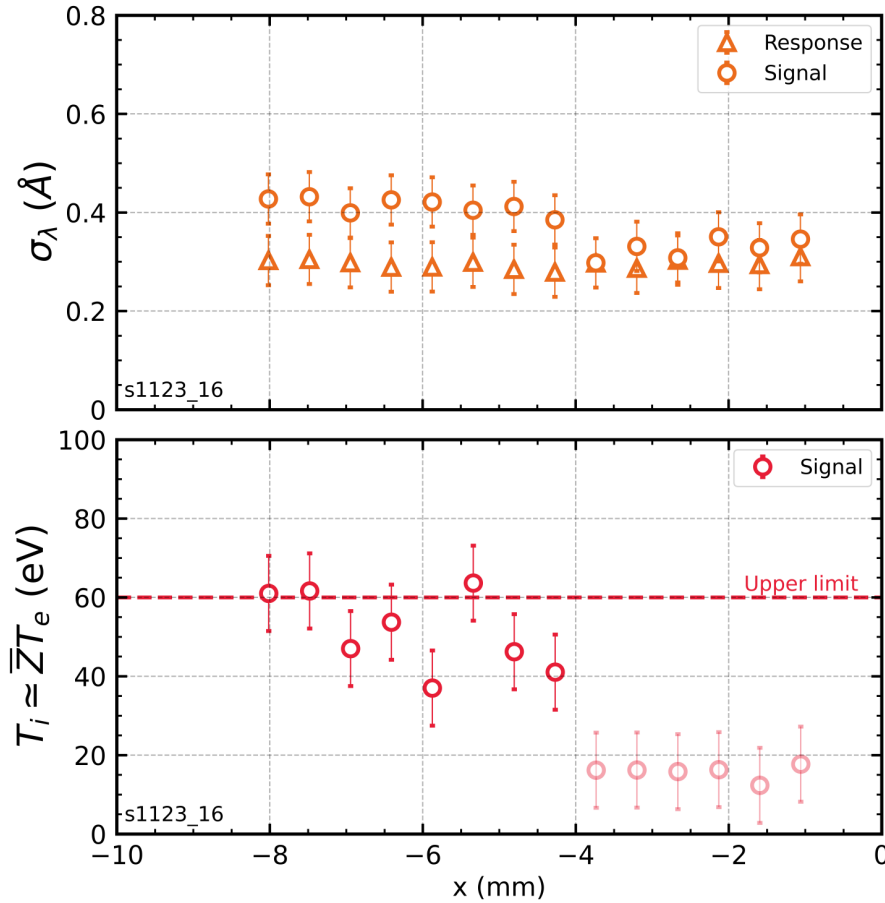


Figure 4.30: (Top) Spatially resolved intensity-weighted standard deviation of the Thomson scattering spectra presented in Fig. 4.27. The error bars are estimated based on the uncertainties of the Thomson scattering spectra, Fig. 4.27. (Bottom) Temperature profile correspondent to the measured spectral broadening, assuming $T_i \approx \bar{Z} T_e$. The shaded data points represent spectra with a signal width equal to the diagnostic response function.

Using the measured spectral width of the Thomson scattering spectra and assuming that the broadening effects are only due to thermal motion, an upper limit of the expected temperature in tungsten plasma can be derived. Fig. 4.30(Bottom) shows the associated temperature to the measured spectral width under the hypothesis of negligible contributions from velocity fluctuations within the scattering volume.

Under these conditions and the absence of visible ion-acoustic peaks, the maximum expected temperature is given by $T_i = \bar{Z}T_e = 60$ eV. A fitting of the spectra is achieved for $T_e \simeq 10$ eV and $\bar{Z} \simeq 6$, where the average charge state $\bar{Z}$ is self-consistently derived using the atomic code, FLYCHK [52]. It is important to note that the absence of ion-acoustic peaks in the Thomson scattering spectra further supports this estimate as, according to the ([84], eq. 5.4.4.), a product $\bar{Z}T_e \simeq 140$ eV is necessary to produce well-separated peaks, much higher than the temperature estimates from the data. In the following section, the same analysis is performed for Thomson scattering data recorded in the experiment in the presence of a planar obstacle.

Flow velocity and spectral broadening measurements in the plasma with the planar obstacle

A similar analysis to the one just presented can be applied to Thomson scattering data taken in the experiment with the planar obstacle. The same Thomson scattering geometry introduced in Fig. 4.25 has been used to perform the measurements. The spectrogram was recorded in the same experiment and at the same time ($t_{TS} = 324$ ns, $t_{shadow} = 354$ ns) as a side-on shadowgraphy image to associate structures in the plasma with the position of the collected spectra. The raw data are shown in Fig. 4.31 together with an example of recorded Thomson scattering spectrum.

Using the similar procedure described in Chapter 2 the location of the scattering volume is determined and overlaid on the shadowgraphy image taken using a 532 nm laser, as shown in Fig. 4.31(a). The Thomson scattering laser vector, k_{in} , is directed along the x-axis, at approximately mid-height of the wire array ($z = -1$ mm), located at $x = -10$ mm. The 14 positions of the scattering volumes are aligned with the laser path and cover a spatial range of $\simeq 8$ mm along the x-axis. The scattering positions are numerically labelled starting from the obstacle (fibre 1) up to the upstream plasma flow near the wire array (fibre 14). At the time of the recorded data, density perturbations

dominate the region ahead of the obstacle until $x = -4$ mm. Half of the available collection volumes fall within the perturbed region, from fibre 1 to 7, while the remaining fibres, from 7 to 14, are located in the upstream flow at $x < -4$ mm.

The raw spectrogram recorded at 324 ns after the current start is shown in Fig. 4.31(b). As observed for the null shot, the Thomson scattering signals are concentrated in a narrow region of $2 \text{ \AA}$ near the 532 nm laser wavelength, as expected for the ion-acoustic feature in the collective Thomson scattering regime. In all the collecting volumes, the spectra of plasma self-emission are small and can be easily subtracted from the data. Overall, the Thomson scattering spectra are recorded in all the available fibres in both upstream and perturbed regions with only spectra from fibres 9, 10 and 14 exhibiting the characteristic double peak feature of ion-acoustic waves.

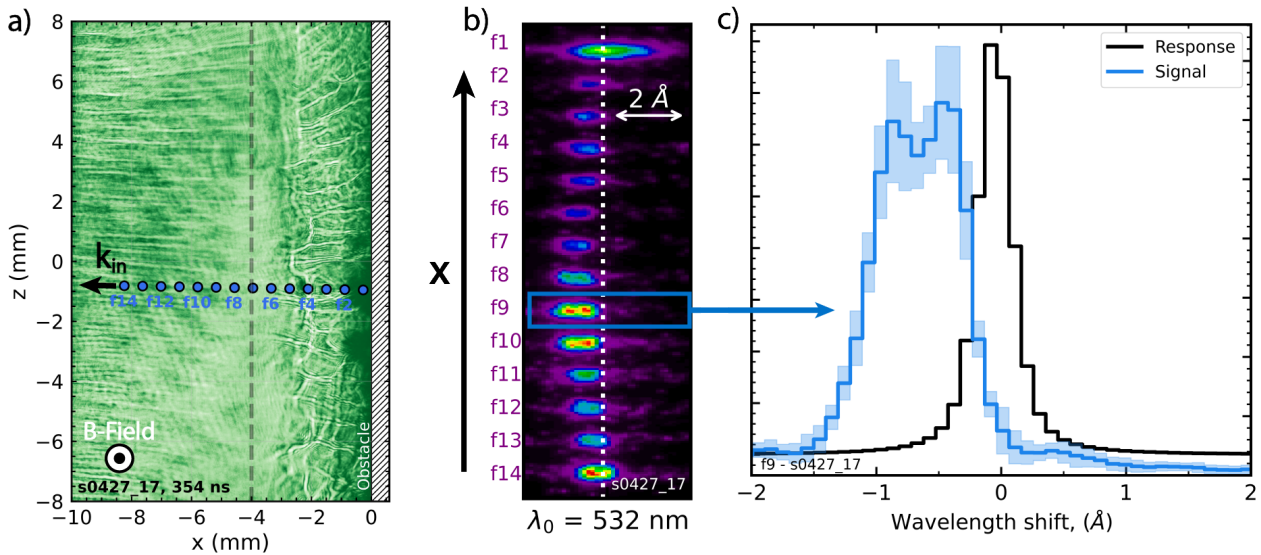


Figure 4.31: (a) Side-on shadowgraphy image of the tungsten plasma flow taken at 354 ns after current start using a 532 nm laser, with the location of the 14 TS collection volumes overlaid. (b) Thomson scattering spectrogram of tungsten plasma flow interacting with a planar obstacle taken 324 ns after current start. (c) Example of Thomson scattering spectrum from fibre 9 ($x = -5.3$ mm). The shaded region represents the standard deviation of the signal, and the solid line is the average over one fibre in the vertical direction of the spectrogram, to boost the signal-to-noise ratio against self-emission and shot noise. The response of the spectrometer to the unshifted laser probe is plotted for comparison.

In the upstream region of unperturbed plasma flow, the collected spectra are visibly Doppler shifted, increasing linearly from fibre 14 to fibre 8. These recorded blue shift signals indicate the direction of the plasma flow from the wire array towards the obstacle. Starting from fibre 7, $x = -4$ mm, the observed Doppler shift decreases smoothly, up to almost zero at the last available fibre in the stagnation

region, $x = -0.5$ mm. This smooth decrease in Doppler shift suggests a smooth flow deceleration within the region $x > -4$ mm which was not detected during the null-shot.

Along with the Doppler shift, significant spectral broadening is detected for spectra collected just ahead of the perturbed layer, $x = -4$ mm and deeper into the stagnation region, $x > -2$ mm. Pronounced ion acoustic peaks are observed in fibres 9 and 10 associated with an increase in spectral width. Spectra where ion acoustic peaks are not observed, fibres 1÷8 and fibres 11÷13, are either regions of cold plasma or the density perturbations observed in Fig. 4.18 give rise to associated velocity fluctuations within the scattering volume which broaden the Thomson scattering spectra and obscure the peaks. Significant broadening effects are also observed at stagnation, from fibre 1 to 4.

An example of a Thomson scattering spectrum collected from fibre 9 is shown in Fig. 4.31(c). The signal is represented by the solid blue line obtained by summing up the intensity over one fibre in the vertical direction of the spectrogram. The shaded area indicates the standard deviation of the signal. The black line represents the response function of the diagnostic recorded from a pre-shot spectrogram as described in chapter 2. In this spectrum, both Doppler shift and spectral broadening are significant compared to the characteristic response. Despite the significance of the broadening signals recorded in this experiment, this effect remains insufficient to reliably measure temperatures from the standard fitting process of the Thomson scattering form factor $S(k, \omega)$, therefore only estimates of the spectral width for each spectrum will be reported.

Following the same analysis described for the free-flow experiment, the velocity component, V_x and the spectral width represented by the weighted average standard deviation σ_λ of the spectra can be derived in the presence of a planar obstacle. The results of the Thomson scattering analysis are summarised in Fig. 4.32 together with the measurements in the absence of obstacle presented previously. From the shadowgraphy image in Fig. 4.32(a) taken at 350 ns after current start, the perturbed region extends between $-4 \text{ mm} < x < 0 \text{ mm}$, while the upstream remains dominated by the periodic density modulation discussed in § 4.3.2.

Fig. 4.32(b) shows the spatially resolved measurements of the flow velocity component, V_x , from the Doppler shift measurements with and without the planar conducting obstacle as described previously and assuming negligible V_y . In the absence of an obstacle (orange circles), the flow velocity

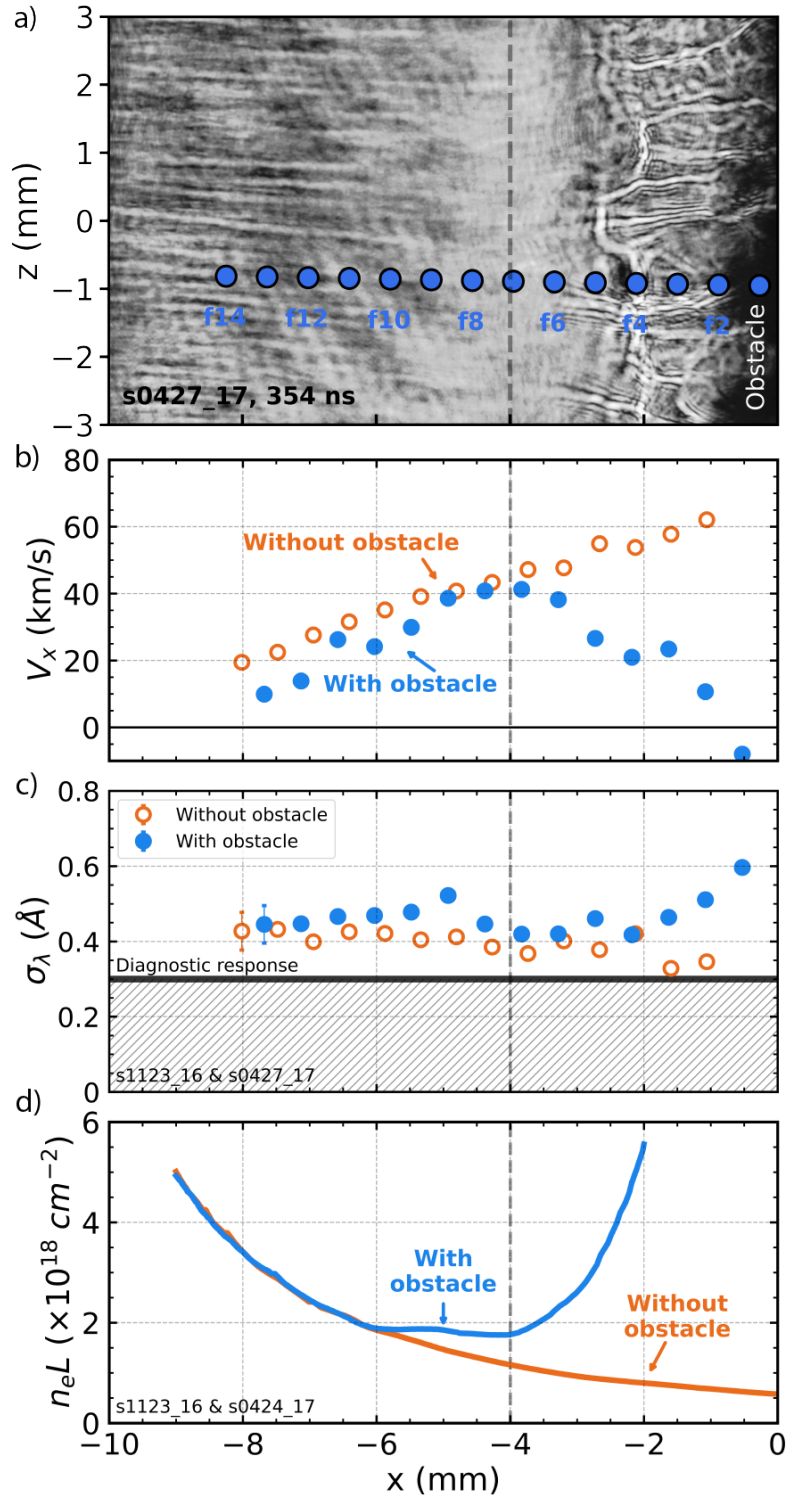


Figure 4.32: (a) Side-on shadowgraphy image of the tungsten plasma flow taken at 354 ns after current start using a 532 nm laser, with the location of the 14 TS collection volumes overlaid. (b) Velocity measurements of the plasma with and without the planar obstacle. The uncertainty in the velocity measurement is negligible (± 2 km/s) resulting in error bars smaller than the size of the markers. (c) Broadening profile of the TS spectra. σ_λ was calculated as one intensity-weighted standard deviation of the spectrum distribution. (d) $n_e L$ profiles taken at $y = 0$ mm from the end-on interferometry data in Fig. 4.23 and Fig. 4.22, at 348 ns after current start with the obstacle (blue line) and 36 ns later (384 ns after current start) without the obstacle (orange line).

shows a roughly linear increase with distance from the wire array, starting from 20 km/s near the wire array at $x = -8$ mm and reaching a maximum velocity of 60 km/s at $x = -1$ mm. When the obstacle is present the flow velocity matches that recorded in the null shot for $x < -4$ mm, reaching 40 km/s. Beyond this point, the velocity profiles for the experiments with and without the obstacle begin to differ. The flow gradually decelerates for the case with the obstacle, from the maximum 40 km/s at $x = -4$ mm to almost 0 km/s at stagnation, $x = -0.5$ mm. This deceleration can be caused by the transverse magnetic field piling up at the obstacle surface, which will be discussed later in the chapter.

Fig. 4.32(c) shows the measured spectral width of the Thomson scattering spectra as weighted average standard deviation, σ_λ . When the obstacle is not present, the scattered signals are only 40% wider than the diagnostic resolution, which is approximately $0.3 \text{ \AA}$, and the width of the spectrum remains nearly constant for all the spatial positions measured. In the presence of the obstacle, the broadening is similar for the first few spatial positions with $x < -6$ mm. However, we observe a measurable increase in $\sigma_\lambda \approx 0.5 \text{ \AA}$ ahead of the perturbed region, at $x = -5$ mm. The spectra become narrower again before becoming much broader close to the obstacle, $x > -2$ mm, in the region which cannot be probed using the interferometer and where significant density perturbations were observed in the side-on shadowgraphy and imaging refractometer. In the experiment with the obstacle, distinct ion acoustic peaks are only visible in fibres 9 and 10, Fig. 4.31.

The absence of the characteristic ion acoustic peaks in most of the spectra and in the free-flow data inhibits our capability to derive accurately the product $\bar{Z}T_e$, and thus $\bar{Z}$ and T_e . This is an issue, especially for high atomic mass plasmas such as tungsten, as broadening due to the splitting of the ion acoustic waves is proportional to $c_s \sim \sqrt{\bar{Z}T_e/m_i}$. Moreover, it is not possible to distinguish the partition between thermal motion and velocity fluctuations within the scattering volumes. This is particularly true in the regions where density perturbations are observed in both self-emission and interferometry data, excluding the hypothesis of uniform velocity at the scale of the scattering volume ($\simeq 150 \text{ }\mu\text{m}$). Based on the spectral width measurements in Fig. 4.31(c) the electron energy required to produce the ion acoustic features is approximately $\bar{Z}T_e \simeq 140 \text{ eV}$. This suggests that a significant increase in the electron energy must occur in these regions compared to the recorded upstream value $\bar{Z}T_e < 60 \text{ eV}$ in this experiment, and in the experiment without an obstacle (fibres 14-11), Fig. 4.30.

The significance of these measurements to characterise the reverse shock structure in tungsten flow will be discussed later in the chapter together with possible explanations of the origin of density perturbations in the region ahead of the planar obstacles.

Magnetic field measurements

Magnetic, or ‘B-dot’, probes are a popular technique for measuring magnetic fields in plasma. In the simplest possible case, a loop of thin wire is immersed in the plasma, and the temporal variation of the magnetic field flux at the probe induces a potential difference across the loop. This potential difference is recorded, and, using the measured area of the loop, the magnetic field can be recovered by integration. Magnetic probes are cheap, straightforward to manufacture and easy to analyse, however, some aspects related to the interpretation of the recorded signals must be considered. It is well known that probes produce plasma perturbations, as they are stationary obstacles which act as thermal sinks and sources of impurities and can be sensitive to temporal variation of the plasma potential. In the tungsten null-shot experiments two pairs of probes have been used to characterise the advected magnetic field frozen in the diverging plasma outflow. The main goal is to determine whether a similar magnetic field to the aluminium case is expected for tungsten and whether this can potentially explain the absence of the stand-off shock.

The design of the probes is shown in Fig. 4.33(b). Two oppositely wound loops are constructed side by side, as close as physically possible to one another. This ensures the probes experience the same magnetic field and electrostatic potential, and the total potential difference around the loops should then be:

$$V_1 = V_p + A \frac{dB}{dt} \quad (4.2)$$

$$V_2 = V_p - A \frac{dB}{dt} \quad (4.3)$$

The probes are coated with insulating tape, which reduces the capacitive coupling of the plasma potential. Additionally, the sign of the potential induced by the magnetic field is opposite on the oppositely wound loops, but the sign of the electrostatic potential is the same in both loops. If the shielding in the loops is perfect, the voltages on the two probes would be entirely due to the magnetic

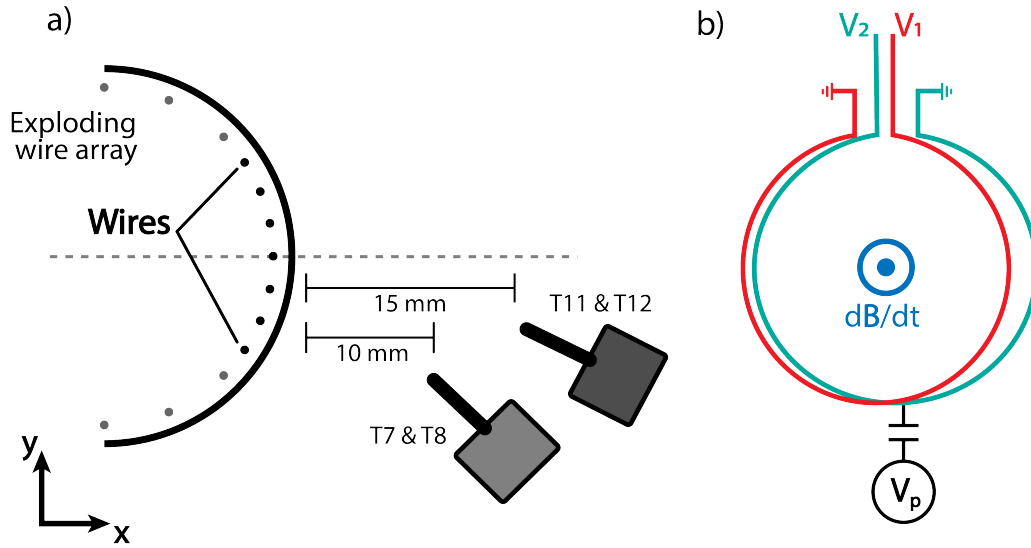


Figure 4.33: Illustration of magnetic field measurements in tungsten free-flow experiment using two pairs of B-dot probes. (a) Locations of the probe pairs in the divergent outflow of a single wire array. T7 & T8 are placed at a 10 mm distance from the wire array (similar distance to the planar obstacle) while T11 & T12 are 5 mm further away, at 15 mm. (b) The probe pairs consist of two loops (shown in red and green), oppositely wound to provide opposite signs for the magnetic potential. The potential across the loops is measured from ground as V_1 and V_2 respectively. The plasma potential, V_p , can be capacitively coupled to the loops.

field and the two probes would show identical signals with opposite signs. When some of the plasma potentials are coupled to the probes, the voltages both show an additional signal with the same sign. In this case, the potential difference due to the magnetic field can be retrieved by taking the difference of the voltages in the two loops, which cancels out the identical electrostatic component:

$$V_1 - V_2 = 2A \frac{dB}{dt} \quad (4.4)$$

The raw signals are shown in grey for the probe pair T7 & T8 in Fig. 4.34(a) - visually, they appear very similar, with opposite signs. This suggests that the electrostatic component must be small. As the electrostatic component is small it is possible to be confident that the signal recorded is mostly due to the voltage induced from the changing magnetic field, which can then be time integrated to give the magnetic field shown in Fig. 4.34(b). From the integrated signal the flow velocity of the plasma can be estimated from the time at which the first signal, and hence the first plasma, arrives at the probe. The time along the horizontal axis in Fig. 4.34 is measured from the current start, and the probe is placed 10 mm away from the array. Using the time of flight as around 80 ns implies a flow velocity of around

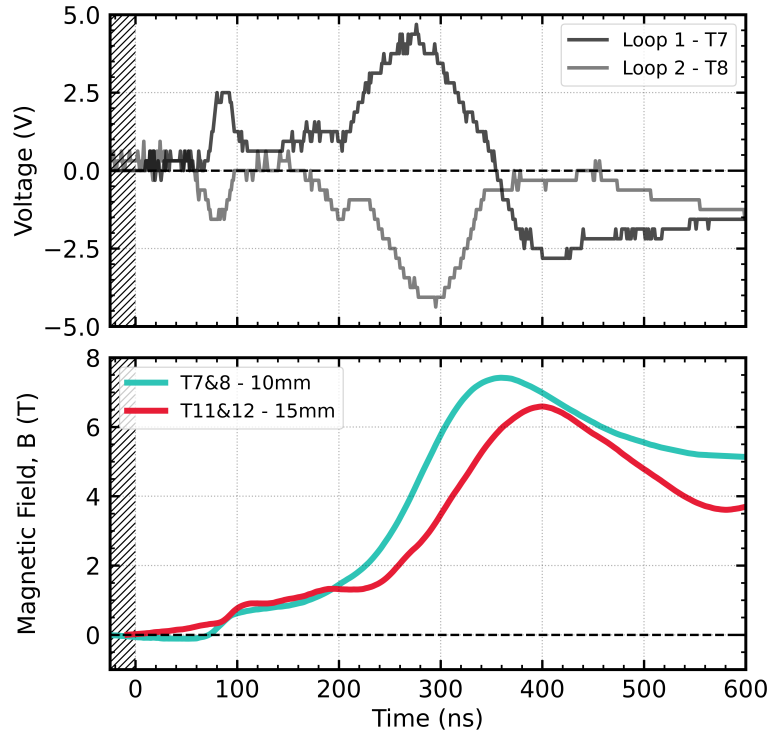


Figure 4.34: (Top) Example of voltage signals recorded by T7&T8 B-dot pairs in the null-shot (s1123_16). The pair of probes were placed in the diverging outflow of the exploding wire array. (Bottom) Measured magnetic field profiles respectively at distances 10 mm and 15 mm from the wire array.

80 km/s which is close to the 60 km/s observed later in time by Thomson Scattering. The maximum recorded magnetic field in both probe pairs is around $\simeq 7$ T at 350 ns and 400 ns after the current start, 10 mm and 15 mm far from the array, respectively. These measurements are consistent with the ones performed in aluminium wires at a comparable distance [59, 60] in a similar experimental setup. Although high-Z elements exhibit a stronger magnetic diffusivity than low-Z plasma since they tend to be colder and Spitzer-Braginskii resistivity is proportional to $\propto T_e^{-3/2}$, this difference appears not significant to reduce the advected field in the plasma as the spatial scale of the plasma is longer than estimated resistive diffusion scale. This implies that a similar magnetic field pile-up is expected in tungsten flow as for aluminium. The average magnetic field accumulated at the planar obstacle can be estimated from the change in magnetic flux in the stagnation region:

$$B_{cum} = \frac{V_{flow}}{x_s} \int B(t) dt \quad (4.5)$$

Where V_{flow} is the incoming flow velocity and x_s is the length of the stagnation region. Assuming

an average flow velocity $\simeq 60$ km/s and a distance $x_s \simeq 2.7$ mm the accumulated magnetic field is $B_{cum} \simeq 4$ T at 300 ns after the current pulse. This value is consistent with previous aluminium experiments, indicating that the transverse magnetic pressure is expected to affect the overall dynamic of the upstream flow. In the next section, the implication of the advected magnetic field on the upstream tungsten plasma will be discussed.

4.4 Discussion of the accretion shock in tungsten flow

Various laser-based diagnostics including shadowgraphy, imaging refractometer and Thomson scattering have been applied to investigate the reverse shock structure formed from the collision of super magneto-sonic plasma flows against a planar obstacle. Data from experiments with aluminium and tungsten exploding wire arrays revealed two completely different behaviours. In the aluminium experiment, a stand-off shock is formed mediated by the pile-up of the transverse magnetic field ahead of the conductive obstacle, which is consistent with earlier work [59, 60]. In contrast, a quasi-one-dimensional shock is not formed in tungsten plasma, instead, small-scale density perturbations are dominant in the interaction region. Overall the latter reverse shock structure can be seen as formed from two distinct regions: a dense stagnated plasma adjacent to the obstacle and an extended low-density region ahead of the stagnation layer. Both are highly non-uniform with an extent similar to the distance of the stand-off shock from the obstacle in aluminium experiments. This reverse shock resembles the complex shock structure discussed in the astrophysical literature [26, 27, 58]. In these cases, the 3-dimensional nature of the system and the formation of density structure can modify the typical shock profile, leading to accretion columns which cannot be described assuming a simple 1-D reverse shock model.

In this section, we discuss the key findings from the data analysis, focusing first on the properties of the shock structure before moving to the discussion of the oscillation along the shock observed from the optical self-emission. A discussion of the small-scale perturbations dominating the upstream flow is provided along with the results from the aluminium experiment. Finally, this chapter concludes with possible physical explanations that could cause the observed reverse shock structure.

4.4.1 Summary of the plasma parameters in tungsten plasma flow

In the previous section, hydrodynamic parameters of the upstream plasma flow in tungsten experiments have been successfully characterised. In particular, flow velocity and electron density were measured using Thomson scattering and laser interferometry in both experiments with and without a planar obstacle. The tungsten plasma flow generated by the inverse wire array was found to be super fast magnetosonic with a flow velocity $V \simeq 40 \text{ km/s}$ and a typical density $n_e \simeq 1 \times 10^{18} \text{ cm}^{-3}$. An upper limit on the temperatures was also estimated with $T_e < 10 \text{ eV}$, $T_i < 10 \text{ eV}$, and an average charge state $\bar{Z} < 6$. For flow-obstacle interaction, the plasma parameters remain the same (unaffected) until $\simeq 6 \text{ mm}$ from the obstacle. Flow velocity then starts to gradually decrease from $V \simeq 40 \text{ km/s}$ to $V \simeq 20 \text{ km/s}$ at $\simeq 2 \text{ mm}$ from the obstacle before dropping to $\simeq 0 \text{ km/s}$ within the stagnation layer. This dense layer of stagnated plasma has a well-defined but non-uniform boundary that occupies a $\simeq 2 \text{ mm}$ layer adjacent to the obstacle surface. A summary of the plasma parameters is presented in Table 4.1 at location $x = -4 \text{ mm}$ in the deceleration region.

In experiments with aluminium plasma, a stand-off shock is observed and flow velocity drops significantly across the shock until it becomes subsonic in the post-shock region. This feature is the result of an achieved pressure balance between the ram pressure of the incoming flow and the thermal pressure in the post-shock region combined with some increase in magnetic pressure. A possible explanation for the observed stand-off shock has been given in [59]. According to early studies, this shock structure could not be explained as a "standard" shock formed by a high Mach number flow, instead, it can be interpreted as a result of a magnetic precursor created by the pile-up of the magnetic flux advected by the flow ahead of the obstacle. Stagnation of the plasma flow containing a frozen-in magnetic flux creates an electric field which drives an electric current on the surface of the obstacle. Changes in the magnetic field in the plasma act differently on the electrons and on the ions. At the spatial scale of the order of the ion inertial length ($\sim c/\omega_{pi}$), ions and electrons decouple and the magnetised electrons are slowed down, creating a magnetic precursor which ultimately leads to the formation of a cross-shock electric field that decelerates the ions.

From Table 4.1, the values of dynamic β defined as the ratio of ram pressure over magnetic pressure is similar between tungsten and aluminium flows with $\beta_{dyn} \gg 1$, however, in experiments with tungsten

Quantity	Symbol	Al	W
Flow velocity (km/s)	V	50	40
Electron density ($\times 10^{18} \text{ cm}^{-3}$)	n_e	4	1
Ion temperature (eV)	T_i	12	< 10
Electron temperature (eV)	T_e	12	< 10
Magnetic field (T)	B	6	4
Average ionization	$\bar{Z}$	3.5	< 6
Mass density ($\mu\text{g/cm}^3$)	ρ	51	45
Ion-ion mean free path (μm)	λ_{ii}	4	5
Ion sound speed (km/s)	c_s	14	6
Alfvén speed (km/s)	V_A	24	17
Sonic Mach Number	M_S	3.6	6.6
Alfvén Mach Number	M_A	2.1	2.4
Reynolds Number	Re	1×10^6	3×10^6
Magnetic Reynolds Number	Re_M	9	3
Dynamic beta	β_{dyn}	9	13
Thermal beta	β_{th}	0.7	0.3

Table 4.1: Typical upstream plasma flow parameters ($x = -5 \text{ mm}$) for aluminium and tungsten inverse wire arrays taken at 350 ns after current start. For aluminium arrays, data is taken from [63] for V [Fig. 3(c)], and [60] for n_e [Fig. 4], T_e and T_i [Fig. 6(c)], B [Fig. 2(b)], and $\bar{Z}$ using the atomic code SpK [85]. For tungsten arrays, the data are taken from end-on interferometry and Thomson scattering measurements presented in Fig. 4.23 and Fig. 4.32. T_e and T_i are inferred from the minimal broadening and $\bar{Z}$ is computed using FLYCHK code [52] assuming $T_e = 10 \text{ eV}$. λ_{ii} is calculated using the thermal velocity of the ions. To calculate Re and Re_M we use $L = 5 \text{ mm}$ as a representative length scale for the pile-up of the magnetic field or the slowing of the bulk flow.

plasmas, no sharp velocity drop or sharp density increase is observed at the start of the interaction region. A summary of the two different reverse shock systems is shown in Fig. 4.35.

Advection of frozen-in magnetic field by plasma flows from the wire array should lead to a pile-up of the magnetic field ahead of the obstacle, thus creating an increase of magnetic pressure. The expected field in the experiments with tungsten plasma flow at the deceleration region is $B_{cum}(300 \text{ ns}) = 4 \text{ T}$, similar to what was measured in previous experiments with aluminium [59], leading to a magnetic pressure $P_B = B^2/2\mu_0 \simeq 6.4 \text{ MPa}$. The ram pressure of the incoming flow calculated at the beginning of the deceleration region ($x = -4 \text{ mm}$) is given by $P_{ram} = n_i m_i V^2 \simeq 82 \text{ MPa}$, hence the dynamic beta, $\beta_{dyn} = P_{ram}/P_B \simeq 13$. Under this condition, the magnetic field can be partially responsible for the deceleration of the flow as we do not observe any increase in temperature from the Thomson scattering measurements shown in Fig. 4.32(c), however, a magnetic field of a factor of ~ 3.5 higher is required to achieve pressure balance. The magnetic Reynolds number defined as the ratio between flow

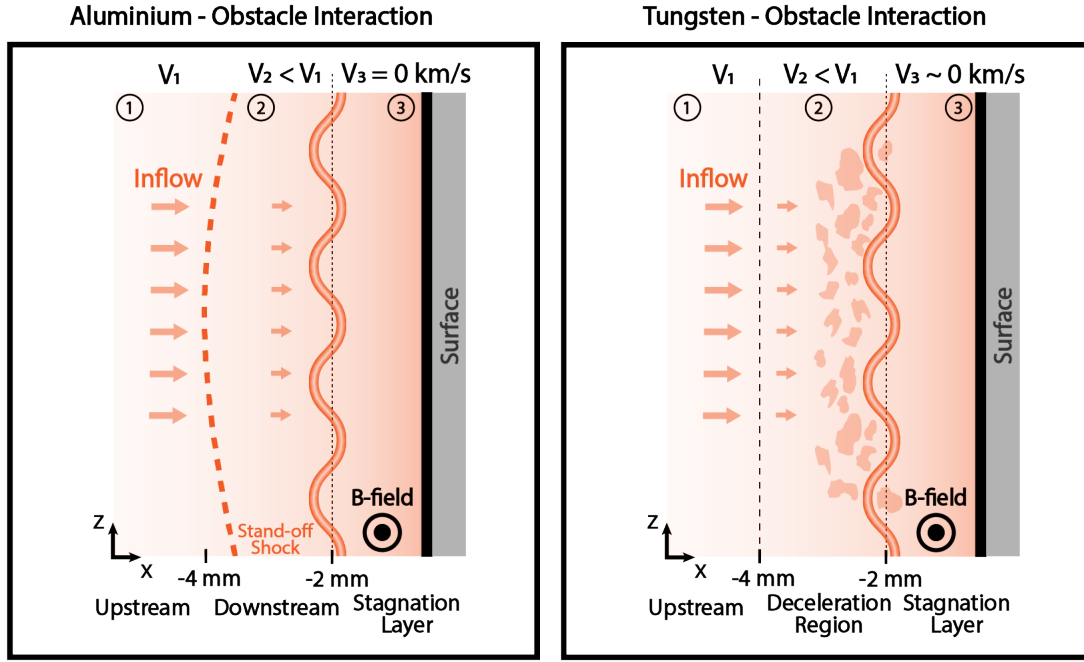


Figure 4.35: Schematic illustration of the two reverse shock structures observed in experiments with aluminium and tungsten plasma flows. (Left) Aluminium-Obstacle interaction. (Right) Tungsten-Obstacle interaction.

advection and magnetic diffusion, $Re_m = VL/D_m$, is also similar between the two materials at spatial scale $L \simeq 5$ mm (see Table 4.1). This means that even in the aluminium-obstacle interaction region, a sharp B-field jump is not expected on the spatial scales of the observed stand-off shock transition ($\simeq \lambda_{ii}$), thus pressure balance should be achieved by thermal pressure.

Overall, for both materials (aluminium and tungsten) magnetic field pile-up is expected on a few millimetres spatial scale, ($Re_M = 1 \rightarrow L_\eta \simeq 3$ mm), so in both cases there is no possibility for the B-field jump to act on a collisional spatial scale. The observed density jump in experiments with aluminium plasma happens at spatial scales much smaller than the characteristic diffusion length, $\delta \ll 0.2$ mm, indicating that the stand-off shock in aluminium plasma is due to a jump in thermal pressure, sufficient to balance the ram pressure of the incoming flow. In tungsten plasma flow, temperature jumps are not observed on spatial scales resolvable by Thomson scattering diagnostic, most probably due to very strong radiative cooling effects and ionisation losses (very small effective polytropic index, $\gamma \simeq 1$ as referred to isothermal shocks in astrophysics) which can lead to the absence of steep variations in thermal pressure and consequently to the absence of a stand-off shock. In addition, the presence of strong radiative cooling combined with the possible presence of collisional heating on very small

spatial scales could be responsible for the seeding of radiative cooling instabilities, explaining the observed density perturbations in tungsten flow as will be discussed later in the section.

Despite dissimilarities between tungsten and aluminium reverse shocks in the magnetic precursor region ($x = -4$ mm, Fig. 4.35), a similar stagnation layer is observed ahead of the obstacle ($x = -2$ mm). This layer of stagnant plasma expands upstream as more material accumulates at the obstacle surface. In the experiments with tungsten plasma, this layer expansion happens at almost double the velocity, $v_s \simeq 1.7$ km/s, than the aluminium case ($v_s \simeq 1$ km/s). This is not surprising as the presence of a stand-off shock limits the amount of plasma flow reaching the stagnation layer in the timescale of the experiment. At late time, the velocity in the post-shock region is subsonic in aluminium-obstacle interaction [region 2, Fig. 4.35(left)] while flow remains supersonic for the entire deceleration region in the tungsten reverse shock [region 2, Fig. 4.35(right)], indicating that the shock transition is not resolved by laser imaging diagnostic and it should happen within the stagnation layer. In both cases, spatial perturbations along the boundary of the stagnation layer are observed, especially in tungsten plasma where a wavelength of $\lambda \simeq 3$ mm is measured from optical self-emission and shadowgraphy images. Details of the internal structure of the aluminium and tungsten stagnation layers are currently unknown due to the lack of direct measurements, however, a possible interpretation for both the stagnation shock and the nature of the perturbation is presented in the Appendix A, instead, the focus of this section will be on the density perturbations observed in the deceleration region of tungsten reverse shock.

4.4.2 Spatial scale and structure of the density perturbations in tungsten flows

In this section, the nature of the small-scale density perturbations dominating the deceleration region ahead of the tungsten shock is discussed, using the measurements from the novel imaging refractometer presented early in the chapter.

Side-on shadowgraphy and imaging refractometer data revealed the presence of density structures in tungsten flow ahead of the stagnation shock, in the deceleration region. The perturbations span over a wide range of spatial scales, resembling turbulence. This is in contrast with the observations

in aluminium plasma, where instead, a quasi-one-dimensional shock was observed mediated by the transverse magnetic field. Evidence of magnetic field pile-up is also observed in tungsten reverse shock, leading to questions on the origin of the density perturbations; Are these density structures generated by perturbations in the magnetic field? or are they driven by other mechanisms, such as radiative cooling, disrupting the formation of a quasi-one-dimensional stand-off shock? Insights on the physical mechanisms underlying these perturbations can be obtained by comparing the spatial scale and distribution of the observed structures with the expected instability theory.

Typically, estimates of the perturbation spatial scale are derived from the shadowgraphy images, applying either a 2-D Fourier analysis in a small sub-region of the perturbation or visually assessing the size of the bright eddies. Both of these approaches, however, are often not precise enough to infer meaningful estimates due to the limited dynamic of the diagnostic and the presence of caustics. Shadowgraphy effects are redistribution of probing laser intensity after propagating through density gradients which results in angular deflections of the beam. This intensity distribution is indirectly connected to the spatial scale of density perturbations and the unfolding process requires a number of assumptions, which often are not applicable when imaging plasmas with large density gradients. The new imaging refractometer overcomes at least some of these limitations, providing a direct measurement of the angular distribution of the deflected beam with a high dynamic range, over 500-fold. The results are measurements directly proportional to the electron density gradients ($\sim \nabla n_e$) rather than the second spatial derivative of the electron density ($\sim d^2 n_e / dx^2$), as described in Chapter 2.

Distributions of laser deflection angles have been successfully recorded using the refractometer technique in the perturbed region of tungsten experiments, as shown in Fig. 4.18. The broadening of these angular distributions is (Fig. 4.19) the result of many deflections of the laser beam due to transverse density gradients, along the probing path. If we consider a single deflection event on a density gradient, the deflection angle $\Delta\phi$ of the ray by an electron density gradient is given by:

$$\Delta\phi = \frac{1}{2} \frac{\nabla n_e}{n_{cr}} \delta y \quad (4.6)$$

where δy is the length scale along the probing direction, ∇n_e is the gradient of the plasma electron

density and n_{cr} is the critical electron density. It is noticeable that eq. 4.6 is simply the discrete version of eq. 3.11, presented in chapter 3, [78]. For a single small "blob" of plasma, the density gradient can be estimated as $\nabla n_e = \Delta n_e / \delta z$ (in the z -direction, perpendicular to the propagation direction of the beam, y), with Δn_e representing the change in electron density in the z -direction over the spatial scale, δ_z . To connect the electron density perturbations to the local measured density, n_e , we can rewrite the electron density change Δn_e , as a fraction $\varepsilon = \Delta n_e / n_e \leq 1$ of the total electron density, giving:

$$\Delta\phi \approx \varepsilon \frac{n_e \delta y}{2n_{cr} \delta z} \quad (4.7)$$

This is the deflection from a single density gradient present along the probing path at length δ_y . As observed from the experimental data, however, the ray will encounter many density gradients as it traverses the plasma due to the presence of small-scale density structures, as illustrated in Fig. 4.36. To progress further, a random walk argument can be assumed.

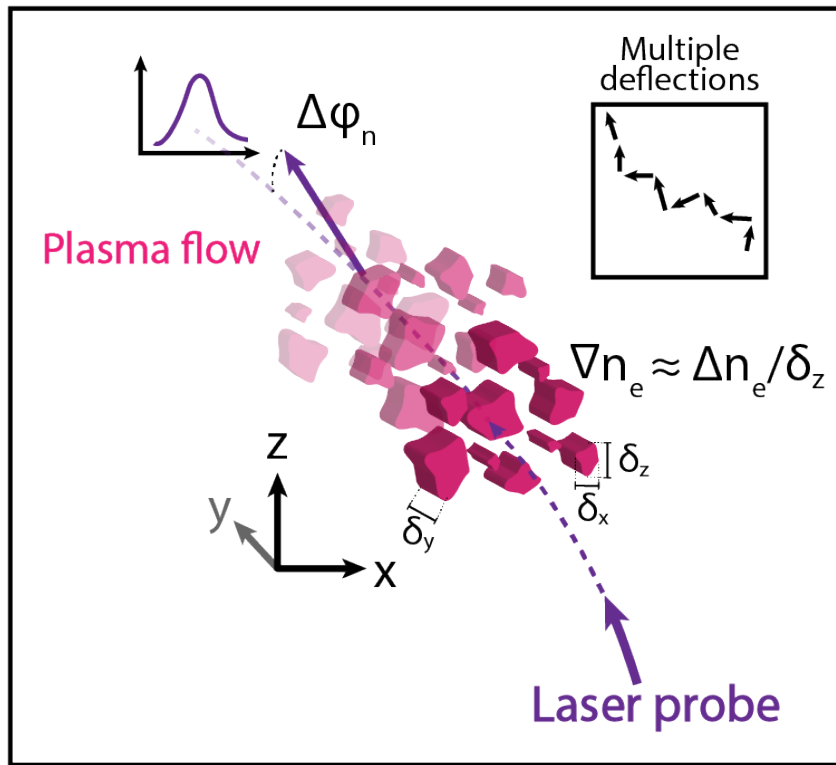


Figure 4.36: Schematic representation of laser probe propagating through a non-homogeneous plasma. The imaging refractometer used in the experiments only measures deflection angles along the z -axis (ϕ), thus the density gradient in z , perpendicular to the probing direction.

Under this assumption, a random walk laser beam propagation through randomly positioned

identical perturbations of density δ_y results in a series of small deflection angles $\Delta\phi$ in the z -direction. After the beam propagates along the entire plasma length, L_y , the probing beam experiences $N_{steps} = L_y/\delta_y$ deflections, thus the total angular broadening of the laser beam is given by:

$$\sigma_\phi^2 = N_{steps} \Delta\phi^2 = \frac{L_y}{\delta_y} \cdot \Delta\phi^2 \quad (4.8)$$

by substituting (4.7) into the random walk relation (4.8), the angular broadening σ_ϕ due to perturbations of amplitude $\varepsilon \leq 1$ can be estimated as follows:

$$\sigma_\phi^2 = \frac{L_y}{\delta_y} \frac{n_e^2 \delta y^2 \varepsilon^2}{4n_{cr}^2 \delta z^2} \quad (4.9)$$

If we assume that on small scales the density perturbations are isotropic so that $\delta y = \delta z$,¹ we can find the spatial scale of density perturbations required to produce a given angular broadening is:

$$\delta_y = L_y \frac{n_e^2}{4n_{cr}^2 \sigma_\phi^2} \varepsilon^2 \quad (4.10)$$

In experiments with tungsten plasma, the measured angular spread from the imaging refractometer data is $\sigma_\phi \approx 3.5$ mrad for the entire perturbed region ($x > -4$ mm, Fig. 4.18), as previously shown in Fig. 4.20. The refractogram was taken using an infra-red 1053 nm laser probe, corresponding to a critical electron density of $n_{cr} \approx 10^{21} \text{ cm}^{-3}$. The average electron density recorded by the end-on interferometry is $n_e \approx 1 \times 10^{18} \text{ cm}^{-3}$, while the total length of the propagation path is $L_y = 10$ mm. Substituting these values in the equation (4.10), we can calculate that the spatial scale for the largest possible perturbation ($\varepsilon = 1$) is approximately $\delta_y \approx 200 \text{ } \mu\text{m}$. If the deflection is produced by smaller perturbations, with smaller amplitude, $\varepsilon < 1$, this would require the presence of structures with spatial scales smaller by a factor of ε^2 to produce the same total deflection σ_ϕ , $\delta_y \approx 200 \cdot \varepsilon^2 (\mu\text{m})$. For example, for $\varepsilon = 0.1$, the spatial scale of perturbations would be $\delta_y \approx 2 \text{ } \mu\text{m}$. To further understand the nature of the observed density perturbations, we can compare the largest estimated perturbation scale with the main characteristic plasma lengths.

¹In principle we can use information from other diagnostics to estimate the ratio of $\delta y/\delta z$ and use this to adjust the estimates which follow. For example, our end-on interferometry in Fig. 4.39 shows $\delta y/\delta z \approx 1/5$. However, this level of detail is not necessary for the purposes of this order-of-magnitude estimate.

Viscous effects in the experiments presented in this thesis are usually negligible and only important at very small scales comparable to the viscous length, $L_v \sim 0.1 \mu\text{m}$. This scale is a few orders of magnitudes smaller than the largest density structures, $\delta \approx 200 \mu\text{m}$, indicating that viscous effects are not responsible for the observed perturbations. An important aspect to consider is that density structures are only obvious from side-on shadowgrams and refractograms while appearing uniform when probed end-on using the interferometry diagnostic [Fig. 4.23(a)]. This cannot be attributed solely to the lower spatial resolution of the interferometry compared to the shadowgraphy diagnostic (due to interference fringes) as density structures are still observed when using side-on interferometry [Fig. 4.24(b)]. Instead, it is more likely that the observed density perturbations are anisotropic, elongated along the y -direction, and randomly distributed in the xz -plane. As a result, they can still be seen from the line integration in the side-on probing direction (y -axis), while in the end-on probing direction (z -axis), they are washed out.

The presence of density perturbations elongated along the magnetic field lines raises the question of what role the magnetic field may play in the pressure balance necessary to support these small-scale density perturbations. Using the inferred plasma parameters presented in Table 4.1 and the Spitzer-Braginskii resistivity, η , the resistive length scale is approximately $L_\eta \approx 1.3 \text{ mm}$. This implies that magnetic diffusion dominates over advection at length scales smaller than L_η , and so the magnetic field is homogeneous on scales below L_η . This resistive length is at least an order of magnitude larger than the perturbation spatial scale estimated from the imaging refractometer data, $\delta \sim 200 \mu\text{m}$, therefore we can conclude that the magnetic field is smooth and not involved in these density perturbations. The unmagnetised nature of the density structures suggests that those density perturbations may rise, instead, due to localised temperature gradients in the plasma.

Thomson scattering and interferometry measurements have revealed changes in velocity and density at scales much larger than the perturbations, however, temperature fluctuations could still occur above the thermal diffusion length, $L_\chi \simeq 150 \mu\text{m}$. This characteristic length is defined as the spatial scale where thermal diffusion is equal to the thermal advection, corresponding to the Péclet number $\text{Pe} = uL_\chi/\chi = 1$, with u the characteristic velocity and χ the electron thermal diffusivity. In this context, thermal energy is mainly carried by the ions due to their larger mass relative to the electrons so the thermal diffusion length is estimated assuming the ion sound speed, $c_s \simeq 6 \text{ km/s}$ and the electron

thermal diffusivity, $\chi = k_{||}/n_e(x = -2 \text{ mm}) \simeq 8.9 \times 10^3 \text{ cm}^2/\text{s}$ (for more details see table 4.2). The estimated value of the thermal diffusion length is comparable to the estimated spatial scale, $L_\chi \sim \delta$, suggesting that thermal conduction may play a significant role in setting the lower limit of the spatial scale of the observed density fluctuations. This aspect combined with the anisotropic nature of the perturbations and the strong radiative properties of tungsten plasma, further motivate us to discuss the role of radiative cooling and its effects on the plasma density.

4.4.3 The role of radiative cooling effects

In the previous section, the main physical lengths have been compared to the estimated spatial scale of the perturbations, revealing that density structures in the tungsten plasma could be supported by localised temperature fluctuations, possibly driven by strong radiative cooling. To further understand the role of cooling effects and the potential rise of instabilities in tungsten plasma, we need first to compare the typical radiative cooling time and length with the corresponding experimental scales in both aluminium and tungsten experiments. In this discussion, we will be utilising the parameters provided in Table 4.1 and the cooling functions obtained from tabulated values calculated using FLYCHK [52].

It is important to note that the cooling function data from FLYCHK are relatively sparse in density and temperature, with intervals of every decade in cm^{-3} for density and every 1 eV in the relevant temperature range. Additionally, due to the complexity of electron energy levels in high atomic number elements such as tungsten, uncertainties in the cooling function persist across atomic codes. Thus, the following analysis should be considered indicative, and to establish a quantitative comparison between experiments and theory, validated high-resolution cooling functions would be required.

As already introduced in Chapter 1, §1.2, radiative cooling effects are relevant to the plasma dynamics when those effects act at shorter timescales than the characteristic timescale of the experiments. In this context, an important parameter is represented by the cooling time, defined as the ratio of thermal energy density to radiated power density, $\tau_{cool} = U_{th}/P_{rad}$. Assuming an ideal gas, the thermal energy density is a known quantity, given by $U_{th} = 3T_e(n_e + n_i)/2$ while we can calculate the radiated power

density as $P_{rad} = n_i n_e \Lambda(n_i, T)$ from the cooling functions presented in Fig. 4.37 for both aluminium and tungsten plasmas. Substituting the respective values of electron temperature and densities (Table 4.1), the expected cooling time for tungsten plasma is $\tau_{cool} \simeq 7$ ns, much smaller than the characteristic hydrodynamic timescale, $\tau_{hydro} = L/c_s \simeq 1 \mu s$. Under this condition, cooling effects can quickly remove heat from the system which could lead to the growth of instabilities as already introduced in Chapter 1 and the most relevant unstable mechanism in this context is represented by the thermal instability.

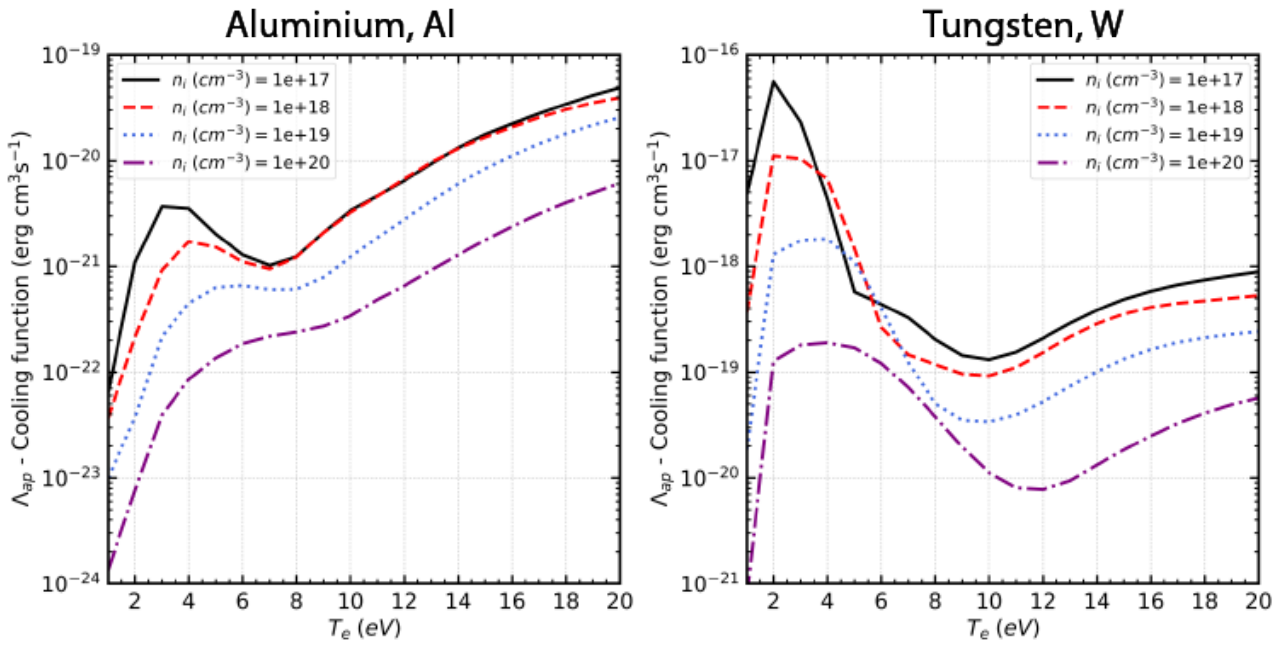


Figure 4.37: Radiative cooling curves were generated using FLYCHK code for both aluminium and tungsten elements. The curves are expressed in normalised power loss ($\text{erg s}^{-1} \text{cm}^3$) obtained by dividing the cooling curves Λ (erg s^{-1}) by the respective ion density, n_i (cm^{-3}).

Thermal instability is a well-known phenomenon in the astrophysical literature, following the analysis of Field [35]. When localized cooling occurs in a plasma, the temperature decreases, resulting in a reduction of thermal pressure. This compression of the cooler region leads to an increase in density and, consequently, the cooling rate. This process becomes a runaway effect, stabilised at small length scales by thermal conduction and at large length scales by hydrodynamic motion. As a result, pockets of dense plasma with various spatial scales dominate the plasma flow, exhibiting a characteristic fragmented density distribution similar to what was observed in tungsten experiments. In this context, the cooling timescale fundamentally determines the growth rate of the instability, and it

is significantly longer for aluminium, $\tau_{cool} \simeq 110$ ns, compared to tungsten, indicating that cooling instabilities could grow much more rapidly in the case of High-Z plasmas. A formal analysis of the thermal instability described in (Field, 1965) yields a cubic equation with three roots representing modes of instability. The stability of these modes depends on the derivatives of the radiated power density with respect to density and temperature. In the limit where the cooling time is much longer than the hydrodynamic time ($\tau_{hydro}/\tau_{cool} \ll 1$ with $\tau_{hydro} = L/C_s$), the cubic equation reduces to a single solution which depends on the derivative of the radiated power density with respect to temperature and thermal diffusivity, $d \log \Lambda / d \log T$.

Fig. 4.38 shows the stability thresholds for different thermal instabilities, isobaric (< 2), isochoric (< 1) and adiabatic (< -0.5). The condition of thermal instabilities is met when the derivative $d \log \Lambda / d \log T$ falls below each threshold. Focusing on the less constrain condition of thermal instability represented by the isobaric case, aluminium plasma appears to be stable for a range of electron temperatures, below $T_e < 9$ eV, while it is predicted to be thermally stable almost independently by the ion density for $T_e > 9$ eV. From previous experiments with the aluminium exploding wire array, an electron temperature between 12 -20 eV has been recorded, following in the stable region predicted by the analysis. Tungsten flow exhibits a similar trend, however, here instabilities are expected for $T_e < 11$ eV and only a narrow stable region emerges from 12 eV up to 15 eV. This again is consistent with the tungsten plasma parameters, where it is measured an upper limit of 10 eV for tungsten electron temperature.

The largest expected spatial scale of the thermal instability is represented by the isobaric length, $\lambda_{iso} \sim C_s \tau_{cool}$, defined as the length scale below which the radiative cooling-induced pressure perturbations can grow. For aluminium, this length is determined to be $\lambda_{iso} = 1400 \mu\text{m}$, while for tungsten, it is significantly shorter at $\lambda_{iso} = 50 \mu\text{m}$. At small scales, the mode can be stabilized by the balancing effect of thermal conduction, preventing excessive radiative loss. This stabilization occurs at a characteristic length scale known as Field's length, expressed as $\lambda_{Field} \sim (\chi \tau_{cool})^{1/2}$ [35].

To calculate the Field's length, it is necessary to determine the thermal diffusivity. Notably, the presence of a magnetic field in the experiments induces an anisotropic thermal conductivity, with reduced electron heat transport perpendicular to the magnetic field. Since the ions in our experiment

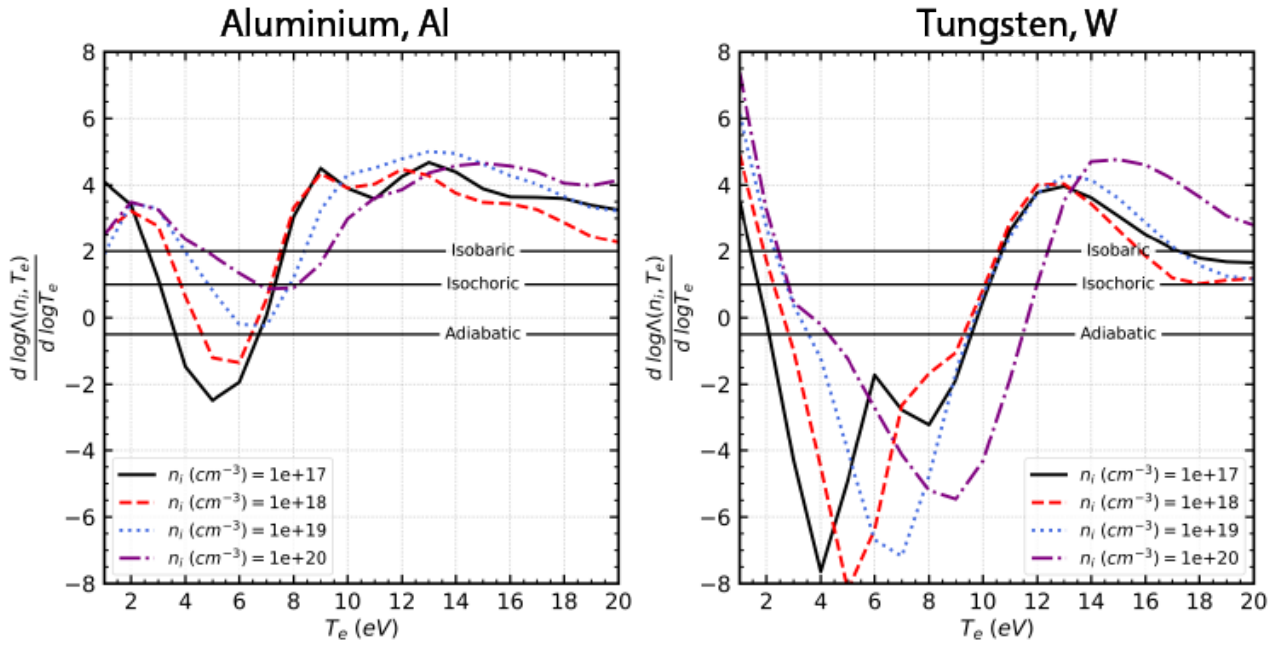


Figure 4.38: Stability analysis of aluminium and tungsten plasmas based on the cooling curves of Fig. 4.37. The horizontal lines indicate various thresholds of instabilities, isobaric, isochoric and adiabatic respectively. Regions above these thresholds are stable regions while modes are allowed to grow when the derivative $d \log \Lambda / d \log T$ falls below the critical values.

are unmagnetized, the ion heat transport is negligible compared to the electron heat transport by a factor of $(m_e/m_i)^{1/2} \sim 10^{-3}$. Consequently, we focus on electron heat transport in the subsequent analysis. By considering representative parameters from Table 4.2, we calculate the Hall parameter $\omega_{ce}\tau_e$ to be approximately 1, where $\omega_{ce} = eB/m_e$ and τ_e represents the electron-ion collision time ([104] Eq. 2.5e). Despite the modest Hall parameter, we find that the dimensionless perpendicular electron thermal heat conduction is three times smaller than the parallel value, employing tabulated data from [105]. The thermal conductivity, $\kappa = \kappa^c k_B n_e T_e \tau_e / m_e$ ([104], Eq. 2.12), and the thermal diffusivity, $\chi = \kappa / n_e$, are then calculated for both the parallel ($\parallel$) and perpendicular ($\perp$) directions. As expected due to the magnetic field, we observe similar values of $\chi_{\parallel}$ and $\chi_{\perp}$ for aluminium and tungsten, with $\chi_{\parallel} > \chi_{\perp}$. Finally, the Field's length is determined representing the length scale of the smallest perturbation for a radiative cooling instability. For aluminium, this length is estimated to be $547 \mu\text{m}$ along the magnetic field and $228 \mu\text{m}$ perpendicular to the field. However, for tungsten, an issue arises. The field's length is derived under the assumption of a large isobaric length, λ_{iso} . Nevertheless, our calculations yield a relatively small λ_{iso} value of $50 \mu\text{m}$, which is smaller than the estimated parallel Field's length of $\lambda_{Field,\parallel} = 130 \mu\text{m}$. Resolving this contradiction would require

Quantity, Symbol (Units)	Al	W
Cooling function, Λ_{ap} (10^{-22} erg cm ³ /s)	9.8	210
Radiated power density, P_{rad} (10^{12} erg/cm ³ s)	270	3900
Thermal energy density, U_{th} (10^6 erg/cm ³)	31	28
Cooling time, τ_{cool} (ns)	110	7
Electron magnetisation, $\omega_{ce}\tau_e$	0.7	0.3
Dimensionless thermal conductivity, $\kappa_{\parallel}^c[\kappa_{\perp}^c]$	7 [1.2]	8.2 [2.8]
Thermal diffusivity, $\chi_{\parallel}[\chi_{\perp}]$ (10^3 cm ² /s)	26.4 [4.6]	67 [22]
Field's length, $\lambda_{Field,\parallel}[\lambda_{Field,\perp}]$ (μ m)	547 [228]	130 [74]
Isobaric length, λ_{iso} (μ m)	1400	50

Table 4.2: Parameters for estimating the relevant time and length scales of the cooling instabilities in the plasma. The relevant plasma parameters are taken from Table 4.1. The cooling functions are taken from FLYCHK [52] using $n_e = 10^{18}$ cm⁻³ and $T_e = 10$ eV for both Al and W, as this was the closest data point in the FLYCHK tables. The dimensionless thermal conductivities ($\kappa_{\parallel}^c, \kappa_{\perp}^c$) are obtained using the tabulated coefficients in [105].

exploring the solutions to the cubic equation without taking the limit of $\omega \ll kc_s$ and determining the smallest length scale at which an unstable mode exists - a new, generalized Field's length. However, due to the uncertainties in the plasma parameters of our experimental data and the limitations in the currently available cooling curves, drawing definitive conclusions is not possible. Nevertheless, it is expected that density perturbations on the order of 100 μ m, elongated by a factor of two along the magnetic field lines due to the suppression of thermal conductivity from partially magnetized electrons, would be present. The obtained perpendicular length scales (δ_x, δ_z) of these perturbations align with the length scales measured using the side-on shadowgraphy and imaging refractometer diagnostics. In the end-on interferometry, the parallel length scale (δ_y) is either in the same order or smaller than the fringe spacing, making it challenging to resolve.

Despite these limitations, a thorough examination of the raw interferograms reveals distinct structures in the post-shock stagnation region, which are highlighted in Fig. 4.39. This figure displays a portion of an interferogram captured at 391 ns after the initiation of the current, using a 532 nm laser. The region of interest is located near the wall and is centred on one of the ablation streams. Notably, the brightness of the interference fringes exhibits small-scale modulation, which we attribute to density gradients resulting from the radiative cooling instability, introducing shadowgraphy effects. To aid visualization, several of these intensity modulations in the zoomed-in image are highlighted, clearly illustrating their approximate dimensions of ~ 100 μ m in width and ~ 500 μ m in length.

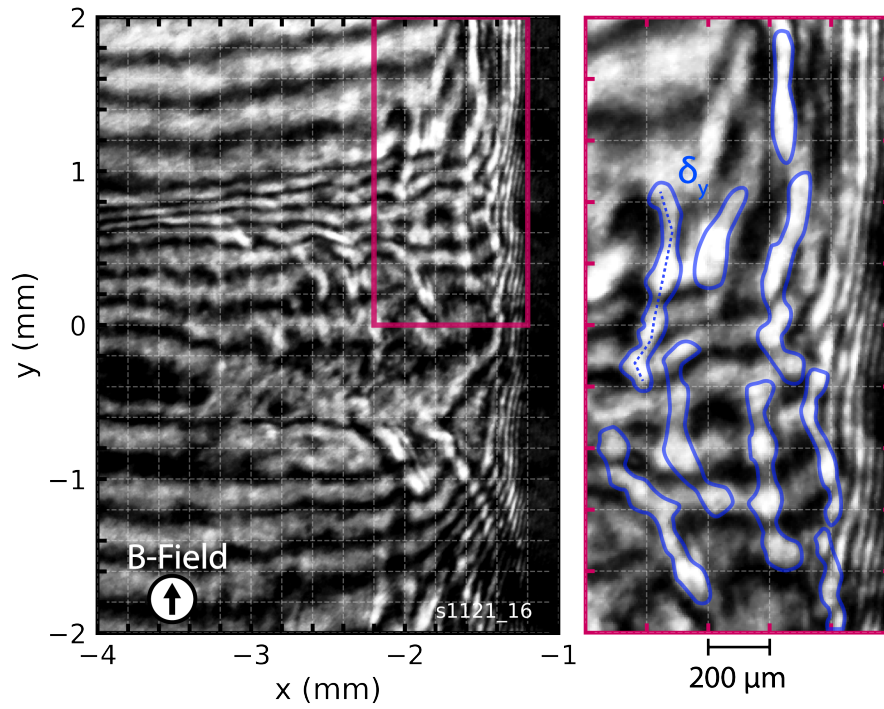


Figure 4.39: A section of a raw end-on interferogram taken at 391 ns after the current start in an experiment with tungsten wires. The section shows a region close to the wall in which the brightness of the interference fringes is strongly modulated, suggesting shadowgraphy effects related to localised structures with large electron density gradients. The zoomed-in region shown on the right is overlaid with blue lines to highlight some of these structures, which are elongated in the y direction (along the magnetic field).

Accurately simulating these experiments requires the use of validated, high-resolution cooling curves that accurately capture the regions where the radiative cooling function decreases with temperature [106]. Consequently, simple free-free radiation models, which only consider bremsstrahlung and recombination, are inadequate to capture the presence of this cooling instability. While the stability analysis presented in Fig. 4.38 agrees with the observed behaviour of both aluminium and tungsten flows at the measured temperatures, it is important to note that the analysis heavily relies on the shape of the cooling functions, which currently lack the required accuracy for detailed predictions and simulations. Therefore, it is not possible to draw definitive conclusions from the threshold and growth rate analysis presented in Fig. 4.38 regarding the thermal cooling instability.

Although radiative effects seem to be more important in tungsten flow than aluminium plasmas, cooling instability could still develop in regions where density is large enough and temperatures are within a specific range. This could be the case of plasma materials in the stagnation layer. From both optical self-emission images (Fig. 4.5) and shadowgraphy data [Fig. 4.7(b)], structures are

observed in the stagnation layer of experiments with aluminium plasma. Here, the growth rate of these perturbations is much slower than in tungsten plasma which is consistent with the estimated long cooling time.

4.4.4 Other possible causes of the density perturbations

Alternative arguments can be made for the possible causes of the density perturbations observed in tungsten experiments. Two possible scenarios that could be relevant and deserve some discussion are electron-thermal instability and kinetic effects.

Electro-thermal instability

The first instability under consideration is the electro-thermal instability (ETI), characterized by a change in resistivity, η , with respect to temperature. Previous studies have observed this instability in gas-puff z-pinchs [107] and discussed it in an attempt to explain the origin of perturbations in metal liners [108, 109] during the implosion. In a plasma where the resistivity follows the Spitzer-Braginskii

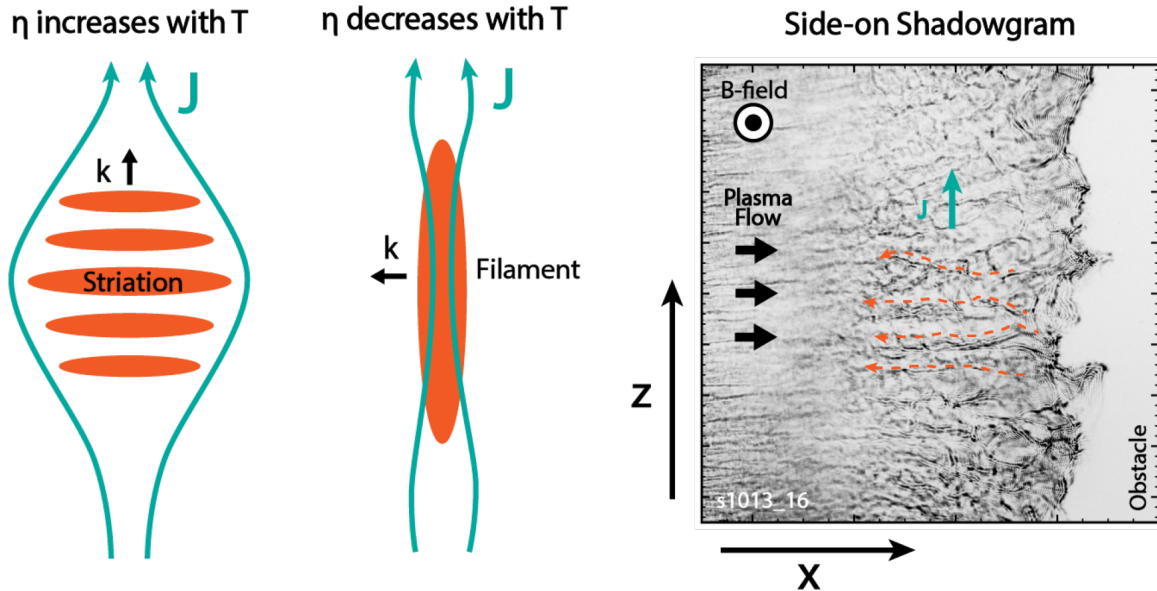


Figure 4.40: Schematic illustration of electro-thermal instability based on the Spitzer-Braginskii relationship between η and electron temperature, T_e , forming striations or filaments. On the right-hand side, a portion of a side-on shadowgraphy from shot s1013_16 taken at $\simeq 440$ ns [presented in Fig. 4.17(b)] (false heat-map for better visualisation).

relationship, $\eta \propto T_e^{-3/2}$, the ETI manifests as filaments aligned with the current flow. Localized Ohmic heating results in a temperature increase, which subsequently decreases the resistivity, facilitating a more favourable path for current flow. This feedback loop amplifies the heating and concentrates the current further. Conversely, in metals where resistivity increases with temperature, localized Ohmic heating elevates the resistivity, causing the current to divert around hotter regions and forming a feedback temperature-heating loop. This diversion leads to perturbations perpendicular to the current direction, forming striations [107].

However, there are two primary reasons why the ETI is unlikely to account for the observed density perturbations in our experiments. Firstly, the length scale of the magnetic field and current density perturbations must be larger than the magnetic diffusive length scale, $L_\eta \approx 1.3$ mm. As a result, the ETI should be ineffective at smaller length scales, preventing the creation of anti-correlated density perturbations on those scales. Secondly, analysis of shadowgraphy data reveals structures predominantly consisting of striations extended along the magnetic field with a k -vector parallel to the current flow direction ($k \parallel z$) rather than filaments with a k -vector perpendicular to it ($k \parallel x$). For the ETI to explain these perturbations, it would require a positive slope of resistivity with temperature ($d\eta/dT_e > 0$), which is characteristic of metals rather than the plasma under investigation. Further investigation of the temperature dependence of the ionization state $\bar{Z}(T_e)$ using FLYCHK ionization tables, [52], indicated that there are no regions where $d\eta/dT_e > 0$. Consequently, the resistivity-temperature relationship in this plasma does not facilitate the formation of ETI striations. The combination of constrained length scales due to magnetic diffusion and the absence of the required resistivity slope leads to the conclusion that the ETI does not explain the observed phenomena.

Kinetic effects

Plasma instabilities can arise from kinetic effects, where the plasma possesses free energy associated with a non-Maxwellian distribution. These instabilities occur as the plasma relaxes towards a Maxwellian distribution. In a previous study by [59] using aluminium wire arrays and a planar conducting obstacle, evidence of reflected ions in the Thomson scattering spectra was observed within 180 ns after the initiation of the current. These reflected ions exhibited velocities comparable to the

bulk plasma flow (100 km/s) but in the opposite direction. It is conceivable that these reflected ions could trigger instabilities such as the two-stream instability, leading to density perturbations. This effect may be more prominent in tungsten plasmas due to the heavier ion mass, resulting in a longer ion-ion mean free path.

However, it is important to note that [59] investigated mostly earlier stages of shock formation compared to our current experiments, when the plasma density was much lower, and thus the ion-ion mean free path was significantly longer. At later times, around 350 ns after the current initiation when our observations were made, Lebedev et al. observed no evidence of reflected ions. Similarly, in our experiments using tungsten wire arrays, we did not detect any evidence of reflected ions in the Thomson scattering data. This observation aligns with the very short ion-ion mean free paths calculated from our experimental parameters, as presented in Table 5.1. Consequently, kinetic effects are thought to not play a substantial role in the formation of the density perturbations observed in our experiments.

4.5 Summary

In this chapter, a series of experiments investigating the formation of a reverse shock formed from the interaction of supersonic magnetised plasma flow with a conductive planar obstacle have been presented. The experiments were conducted using an exploding wire array as a plasma source, extending the past research on aluminium plasma flow to the investigation of reverse shocks in high-Z plasmas, such as tungsten, where radiative cooling effects are expected to be stronger. This is a novel aspect that has not been previously studied with the key findings summarised below:

- The collision of a supersonic magnetised plasma flow against a planar conducting obstacle leads to the formation of a reverse shock, accreting as more material accumulates at the obstacle surface. These shock structures appear very different between experiments with aluminium and tungsten plasma flows despite the same upstream plasma parameters (density ρ , velocity v and magnetic field B). In aluminium plasma, a stand-off shock is formed supported by the pile-up of the magnetic field, extending up to 4 mm ahead of the obstacle. In experiments with tungsten plasma, there is no formation of a stand-off shock, instead, density perturbations at various

spatial scales dominate the region ahead of the obstacle. This region occupies the same spatial extent as the post-shock region in aluminium plasma with a characteristic length comparable to the magnetic diffusion length of the plasma, $L_\eta \simeq 3$ mm. In both experiments, a stagnation layer of dense plasma is also observed near the obstacle surface.

- Measurements of spatial scales of the density perturbations in tungsten plasma flow were performed using the novel imaging refractometer technique. The characteristic size of these density perturbations was estimated from direct measurements of the laser deflection angles along z in the perturbed region, $\sigma_\phi \simeq 3.5$ mrad, assuming random walk and isotropic perturbations ($\delta_y = \delta_z$). The largest spatial scale of the density perturbation was found to be $\delta_y \leq 200 \mu\text{m}$ in the direction parallel to the magnetic field, while transverse to the magnetic field, almost three times smaller due to partially magnetised electrons. This observation is consistent with the expected changes in thermal conduction along and perpendicular to the magnetic field lines.
- An analysis of the importance of radiative cooling effects was performed for both tungsten and aluminium plasmas. The cooling time in tungsten flow, $\tau_{cool} \simeq 7$ ns, is much shorter than aluminium ($\simeq 110$ ns), indicating that the rise of density perturbations in tungsten reverse shock could be associated with strong cooling effects in tungsten plasma, specifically due to thermal instability. This could also explain the absence of a stand-off shock in experiments with tungsten plasma, as heating losses due to radiative cooling can prevent the reach of the necessary pressure balance, while the flow decelerates due to the magnetic field pile-up.
- An additional feature in the reverse shocks is the formation of a stagnation layer near the planar obstacle. As more plasma material accumulates at the obstacle surface, the stagnation region accretes, resulting in a thin and dense structure observed only in the optical self-emission images. This thin-shell layer exhibits corrugations with a characteristic wavelength, $\lambda_0 \sim 3$ mm. Notably, in the case of aluminium, although a similar stagnation layer is observed, the wavelength is significantly smaller and less defined ($\lambda_0 < 1$ mm).

In summary, possible physical mechanisms of the observed reverse shocks in both aluminium and tungsten plasmas have been identified. Although these accretion shock experiments can still be relevant

to astrophysical systems, the presence of a relatively weak transverse magnetic field has a strong effect on the reverse shock morphology, despite a high dynamic β of the plasma flow ($P_{ram} \geq 10P_B$). Here, we have shown that to accurately reproduce astrophysical accretion shocks in the lab, it is necessary to develop a new platform where the magnetic field is parallel to the flow. In addition, it would be beneficial to avoid the use of a solid obstacle, which could disturb the structure of the magnetic field at the stagnation layer and possibly introduce unwanted low-dense plasmas from the ablation of the obstacle surface at the early phase of the experiments. While it is still in the early stages of development, in the next chapter a new platform based on x-ray-driven ablating targets will be introduced, yielding a few interesting preliminary results on colliding plasma flow systems. This platform shows promise in addressing some of the limitations mentioned above and provides potential avenues for future investigations.

Chapter 5

Colliding Supersonic Plasma Flows

This chapter introduces a novel experimental platform specifically designed to investigate stagnation layers that result from the collision of two counter-streaming plasma jets in the presence of an ambient magnetic field. The primary objective of this study was to replicate experimental conditions resembling accretion shocks observed in young star formation and explore the impact of the magnetic field on the stability of shock structures, surpassing the limitations of conventional inverse wire arrays. The radiative platform offers several advantages over traditional inverse arrays. In addition to producing supersonic flows in a highly controllable environment, it allows for self-interacting plasma flow experiments under various magnetic field configurations, thus unlocking new possibilities for studying the role of magnetic fields in many astrophysical systems such as young stars, protostellar jets or Harbig-Haro jets. Accretion shocks are important phenomena that occur when plasma material collides with the chromosphere of a star. The presence of a magnetic field aligned parallel to the flow direction introduces intricate interactions and phenomena that cannot be directly resolved with conventional measurement techniques alone, justifying the need to replicate similar systems in the laboratory.

The experimental platform generates boundary-free, supersonic plasma flows by ablating small solid targets placed parallel to each other on a 3D printed support, see Fig. 5.1. Previous experiments conducted by [67] on single targets, demonstrated that these plasma flows are supersonic and investigated their expansion in a vacuum against the ambient magnetic field of a Z-pinch. By modifying this setup to allow the interaction of two plasma flows propagating along the magnetic field directions,

shock structures form, which can be experimentally analyzed using laser probe diagnostics, enabling a detailed examination of the underlying physics. Importantly, shock structures are not affected by the presence of solid obstacles or physical boundaries often used to collimate the plasma flow (Chapter 1, §1.3). To comprehensively characterize the properties of the plasma in both pre-shock and post-shock conditions, we conducted detailed measurements utilizing laser imaging interferometry, optical self-emission imaging, and Thomson scattering. These measurements provided valuable insights into the main plasma parameters during the shock transitions, revealing the underlying physics governing the structure of the shocked region.

The main content of the chapter is on colliding plasma flows in an ambient magnetic field parallel to the flow direction. It provides a detailed description of the experimental setup and the diagnostic techniques employed. Furthermore, it presents the experimental results obtained from time-resolved optical self-emission imaging, laser interferometry, and spatially-resolved Thomson scattering measurements conducted across the formed shock structure. The findings are accompanied by an extensive discussion of the shock system, including possible interpretations of the observed phenomena. In addition, the chapter introduces a preliminary exploration of colliding flows within a transverse ambient magnetic field, emphasizing its significance in understanding the influence of magnetic fields on colliding outflows and related astrophysical phenomena.

5.1 Collision of two counter-streaming plasma flows in a parallel ambient B-field

The experiments presented in this study utilise a setup similar to the one introduced by [67] and discussed in detail in Chapter 2, §2.3.2, with two significant differences. Firstly, a second source of plasma flow is included to investigate the interaction of two counter-streaming colliding flows. This novel addition raises design questions regarding target separation, size, and material, as it aims to generate a shocked layer, an aspect that has not been previously explored. Secondly, the study examines two different orientations of the ambient magnetic field relative to the plasma flows to understand their effects on the streaming plasma flows. A schematic representation of the experiment and the diagnostic

field-of-view is provided in Fig. 5.1.

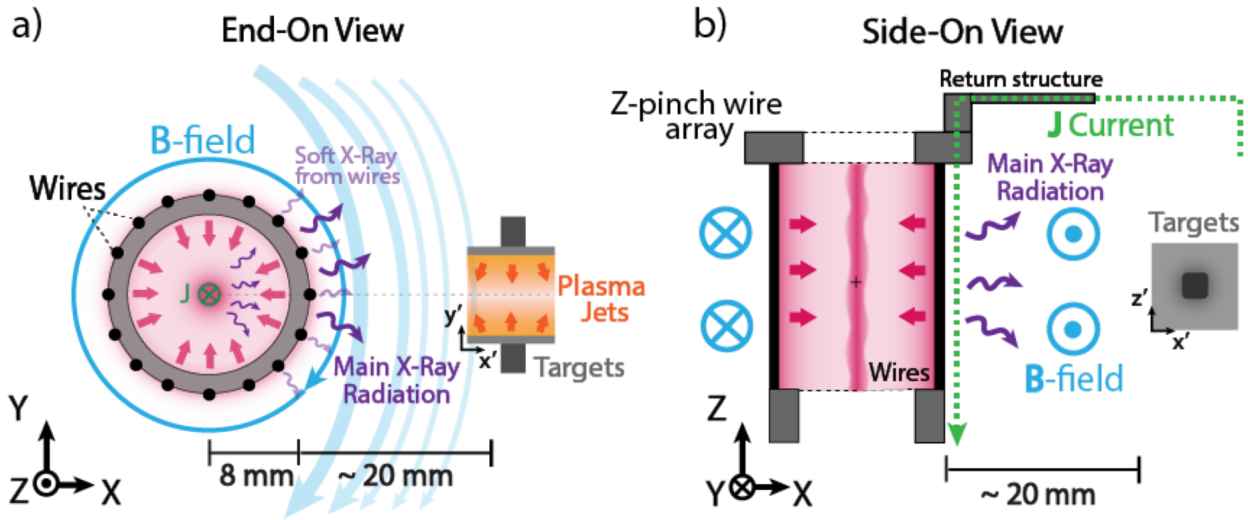


Figure 5.1: Schematic representation of the X-Ray driven colliding flows experiment. (a) An end-on view (xy -plane) of the imploding wire array on the left-hand side and the two parallel targets system on the right-hand side. There are two distinct sources of X-rays, soft low-energy X-ray radiation due to wire ablation at the start of the current pulse and main high-energy X-ray radiation, $\simeq 15$ kJ, produced at full array implosion. The target holder is positioned at $\simeq 28$ mm from the central line of the Z-pinch. The ambient magnetic field is parallel to the plasma jets direction, along the y -axis. (b) A side-on view (xz -plane) of the experimental setup, the ambient field is supported by the current pulse J , passing through wires.

When the current J passes through the return structure and subsequently through the wires, it induces the formation of an ambient magnetic field surrounding the wire array and azimuthally directed along the y -axis at the target region, Fig. 5.1(a)&(b). As the current pulse passes through the wires, the presence of the $J \times B$ force leads to the implosion of the wire array toward its axis which results in the generation of 15 kJ of X-ray radiation pulse emitted symmetrically from the array. This radiation pulse is responsible for driving the plasma flows through ablation from the two solid targets. Targets are positioned in the magnetic field produced by the current flowing through the wire array Z-pinch.

Silicon targets have been chosen as plasma sources for several reasons. Firstly, the outflow of the ablated plasmas from silicon has been extensively studied in single-target radiatively driven experiments, demonstrating its supersonic nature and plasma properties, which are ideal for investigating plasma colliding jets. The profile of the plasma flows is uniform, without any unwanted perturbations that could impact the shock structure and complicate the interpretation of results. Secondly, silicon plasma exhibits behaviour similar to aluminium-based plasmas, which have been thoroughly characterized at

MAGPIE in the context of radiative shock experiments [74]. Furthermore, well-established theoretical and empirical models exist for low-Z elements like silicon and aluminium. Ionization models are available for the plasma parameters presented in the thesis, enabling the inference of average ionization, $\bar{Z}$, and electron temperature, T_e from the measured $\bar{Z}T_e$. The radiative cooling rate can also be estimated by comparing the available cooling curves with the data from [110]. Another very important advantage is that silicon is non-conductive, so no significant magnetic field perturbations by silicon targets are expected when it is placed to intercept the magnetic field flux as in Fig. 5.1. Finally, silicon-based plasmas are well diagnosable, with properties such as temperatures and densities, compatible with the currently available diagnostic suite, enabling high-quality Thomson scattering measurements of the interaction region.

Fig. 5.2 illustrates a more detailed schematic of the experimental platform. The cathode structure, consisting of the cathode plate and the cathode disk, is mounted on the central conductor of the MITL (Magnetic Insulated Transmission Line) using a hollow cup-shaped support. To enable end-on probing of the targets along the z-axis, the wire array Z-pinch axis is offset from the chamber centre by 30 mm. Holes on the cathode plate, top, and side of the support facilitate the direction of the end-on laser probe into the imaging optics via a 45° mirror. An anode disk is attached to the return structure at a height of 40 mm, and the wires are hung from it, making contact with the cathode disk below. Grooves on both disks secure the wires in place. The anode structure consists of four return posts connected to the anode plate, which is then grounded by eight return posts to the chamber breadboard.

The experiments involve the use of two square-shaped silicon targets measuring 5 mm x 5 mm with a thickness of 450 μm (In s0201.22 shot the targets were 10 mm x 5 mm). These targets are precisely positioned at the ends of the 3D-printed target holder arms, with their flat surfaces facing each other. The distance between the two targets, approximately 15 mm (measured in each experiment), is determined by the geometry of the target holder, whose design ensures that the interaction region at the mid-plane is always positioned to correspond to one of the laser diagnostic views. In the case of the magnetic field parallel to the flow directions, the collision of the flows is imaged using end-on view laser probes, while the side-on diagnostics are used to observe the level of plasmas expansion in the orthogonal direction after the collision. The entire target holder is not connected to the array hardware and is secured in place by a series of high inductance posts and a kinematic mount grounded to the

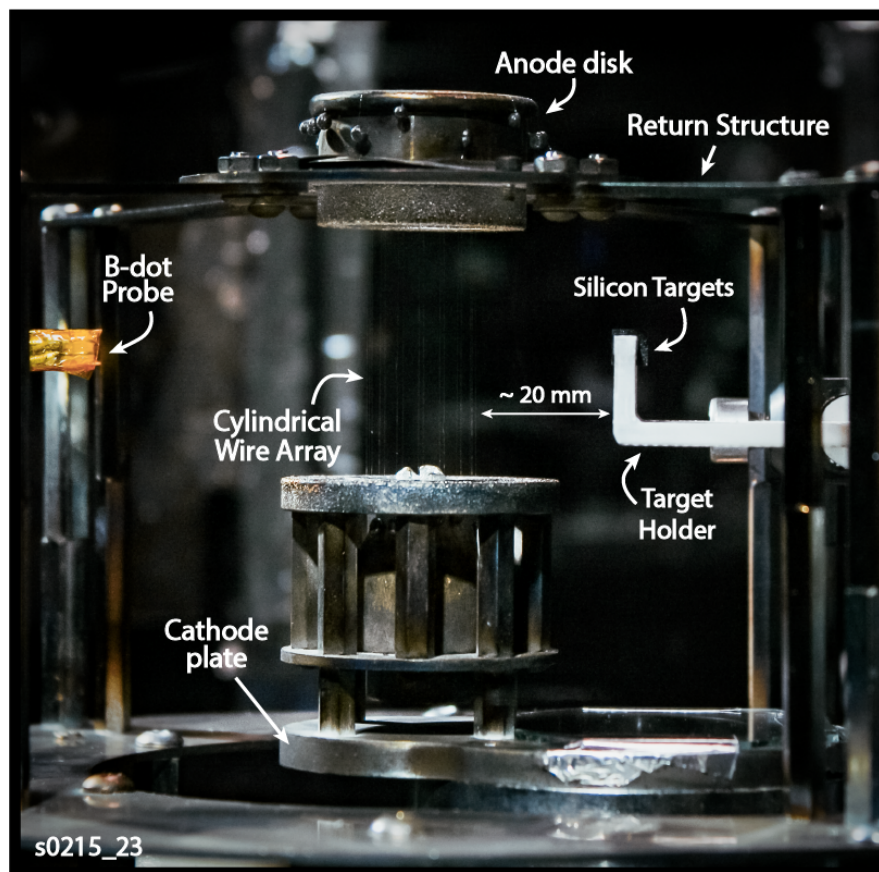
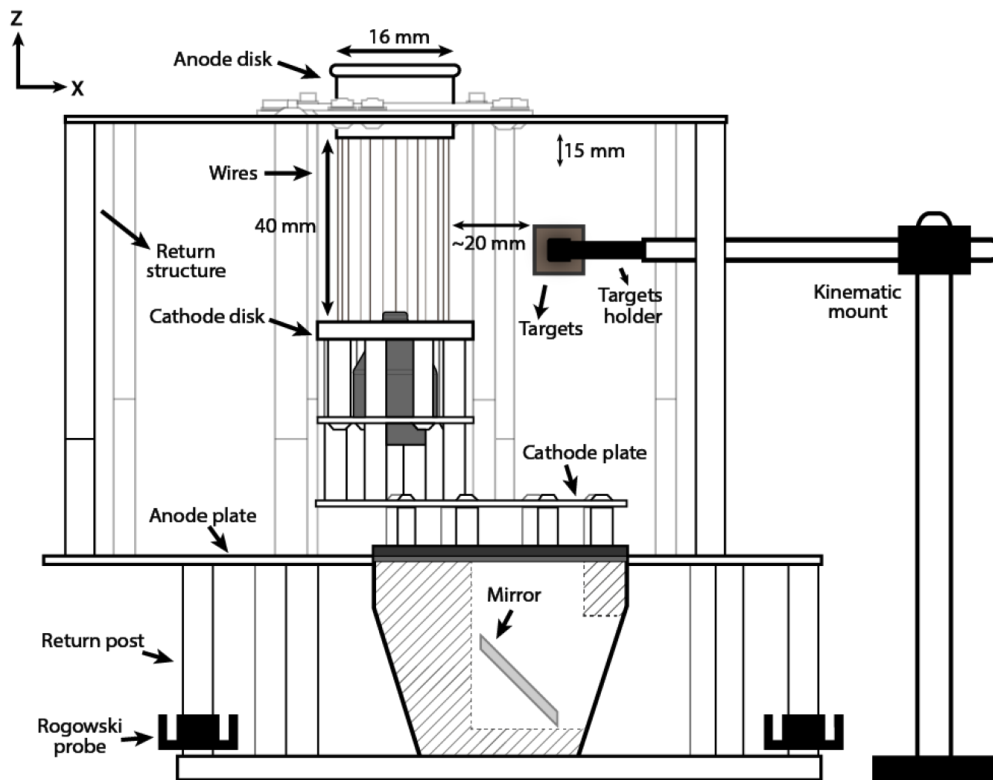


Figure 5.2: (Top) Illustration of the experimental platform lateral view (xz-plane). (Bottom) Photograph of the experimental X-Ray driven platform used for the colliding flows experiment.

chamber breadboard. This prevents direct electric current discharge through the target holder, which could introduce additional circuit paths and plasma ablation from the target holder surfaces. The centres of two targets are positioned at $\simeq 30$ mm from the central axis of the array, with their surface normals oriented along the y-axis, hence along the local B-field direction. Sufficient distance is maintained between the targets and the return current structure to avoid undesired breakdowns. The diagnostic setup for this particular experiment included laser-probe interferometry, optical self-emission imaging, and Thomson scattering. Additionally, conventional machine diagnostics were employed to assess the performance of the MAGPIE generator. Two pairs of Rogowski probes were positioned in opposite return posts, as depicted in Fig. 5.2. The probes in each pair were mounted with opposite polarity, enabling the separate removal of electrostatic pickup from each post and facilitating two measurements of the load current. Figure 5.3 (Left, top) illustrates an example of recorded time-resolved voltage signals from a Rogowski probe during shot s0223_23. The rise in voltage determines the starting time of the current pulse for each experiment. Integrating and scaling the voltage signals yield the reconstructed current profiles, as shown in Fig. 5.3 (Left, bottom), for all experiments discussed in this chapter.

In contrast to plasma flows driven directly by current, the plasma flow formation in experiments with the radiatively driven platform is determined by the X-ray radiation pulse from the pinch rather than the load current. Hence, four filtered photoconductive detectors (PCDs) are employed to measure the X-ray pulse generated by the implosion of the Z-pinch [92, 94]. Fig. 5.3 (Right) presents signals for different filters (Beryllium - $20\ \mu\text{m}$, Polycarbonate - $6\ \mu\text{m}$, Mylar - $2\ \mu\text{m}$ and Aluminium - $6.5\ \mu\text{m}$) arranged from top to bottom. The use of various materials and filter thicknesses allows the detection of a broad range of energy within the X-ray spectrum, enabling detailed characterization of the X-ray pulse. For instance, with a Beryllium ($20\ \mu\text{m}$) filter, the diagnostic becomes sensitive to X-rays above approximately 700 eV, which are only generated after the complete implosion of the array. A Mylar-filtered PCD allows to measure X-ray radiations from plasma precursor, which can be seen as a small signal as shown in Fig 5.3(Right). It is seen that this signal appears 50 ns earlier than the highly energetic X-ray pulse (in this case, it saturates the detector). The X-ray precursor is a result of wire ablation and stagnation of low-dense and fast plasma flows on the array axis, before the main implosion [92–94]. However, this early X-ray emission is not sufficiently energetic to drive dense

silicon plasma flows. The majority of energy driving the outflows originates from the main implosion of the array, making it the most reliable reference time for the dynamics of plasma flows. For more details refer to Chapter 2.

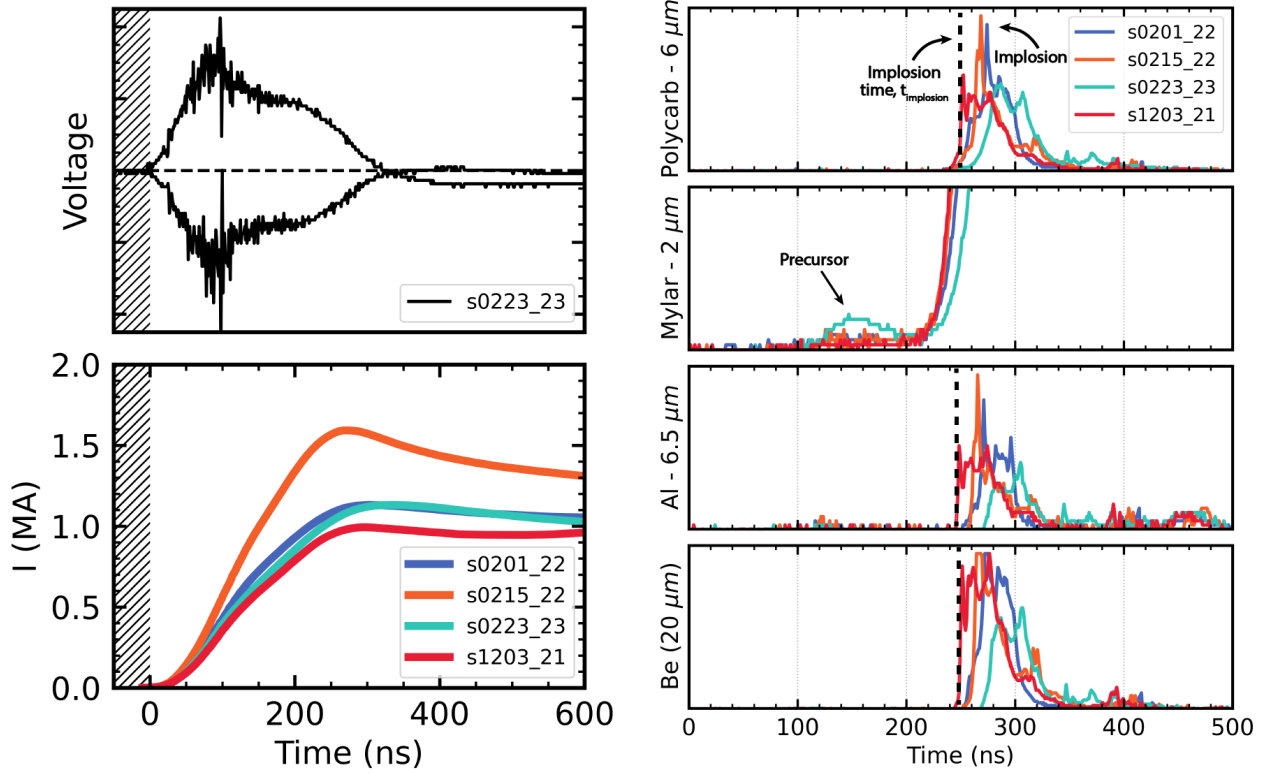


Figure 5.3: (Left) The top panel represents an example voltage signal recorded with Ragowski coils during s0223_23 shot, indicating negligible electrostatic pickup due to signal symmetry. The bottom panel shows the reconstructed current profiles of s0215_22, s0201_22, s1203_21, and s0223_23 shots. (Right) X-ray signals captured using four filtered PCDs directed at the cylindrical imploding wire array. The referenced time for the experiments is from the main implosion: 262 ns after current start for s0215_22 and s0201_22, 270 ns for s1203_21, and 278 ns for s0223_23. The Myler-filtered PCD recorded the x-ray pulse precursor in s0223_23, with the signal saturating at implosion.

By combining the knowledge of the current profile and the initiation of X-ray emission, the timing of plasma formation and collision of the flows can be determined and diagnostic timings synchronized between various experiments. Using the X-ray pulse as a time reference enhances the reliability of comparisons across different shots. For example, the X-ray pulse starts at $t_{implosion} = 262$ ns after the load current for both s0215_22 and s0201_22, while in s0223_23 it starts 16 ns later ($t_{implosion} = 278$ ns). All diagnostics timings recorded for each experiment are synchronized based on the start of the current pulse, t_I , thus a more relevant time for radiatively driven experiments is counted from the start of the X-ray pulse, calculated as $t_{exp} = t_I - t_{implosion}$. In discussing the results, we will always

refer to the timing of particular features based on the time after the X-ray pulse starts unless specified otherwise. For completeness, we will also indicate (shown in brackets) the equivalent time measured from the current start. With the exception of optical self-emission imaging, laser-based diagnostics primarily focus on a narrow temporal range, approximately around 200 ns after the X-ray pulse.

5.1.1 Overview of the plasma flows dynamic from optical self-emission images

A temporal sequence of optical self-emission images recorded using a multi-frame fast camera allows to obtain information about the overall evolution of the plasma system from a single experiment. As already seen in the previous Chapter, the MAGPIE optical self-emission camera is capable of capturing 12 images per shot, with an adjustable frame rate and a 5 ns exposure time. A more detailed description of the optical self-emission imaging arrangement can be found in Chapter 2 § 2.2.2. For

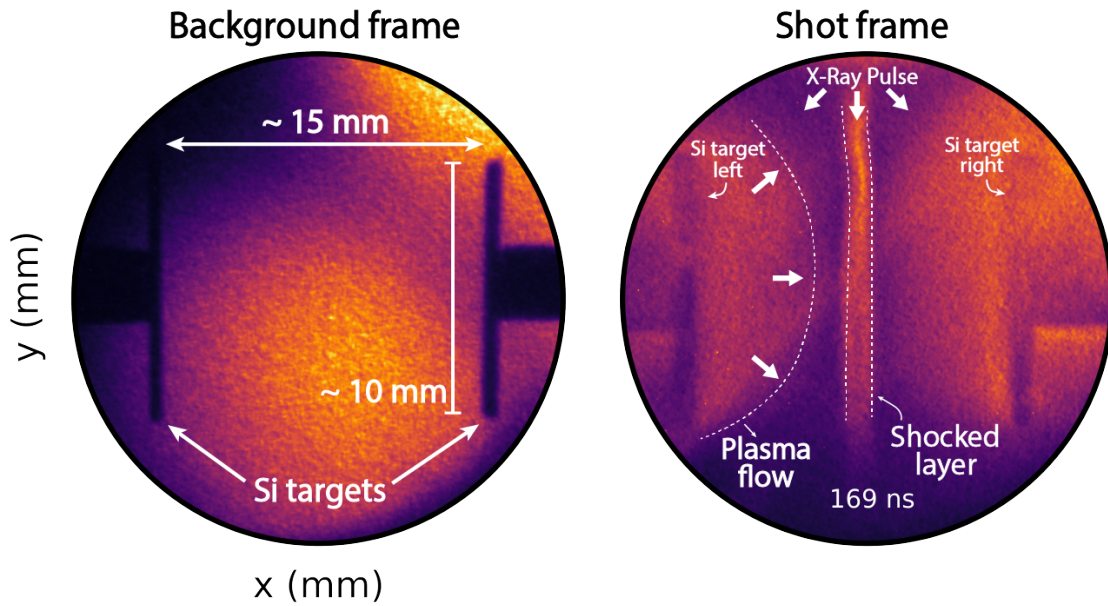


Figure 5.4: Example of optical self-emission background and shot frames from s0201_22 shot taken from the end-on view (xy-plane). The silhouettes of the two silicon targets are visible in both images. The distance between the two targets is approximately 15 mm while the target length is 10 mm. At 169 ns (431 ns) after the X-Ray pulse plasma flows from both targets are visible, propagating mainly along the target surface normals. The shocked layer is formed at the mid-distance between the two targets.

all the experiments presented in this chapter, the temporal range captured by the camera is 300 ns (inter-frame, $\Delta t = 25$ ns), extending from 90 ns after the wire array implosion phase (start of the X-ray pulse) to 390 ns. In this particular camera configuration, the line-of-sight is oriented in the end-on

view, following the same optical path as the end-on interferometry. The camera is focused at the target centre ($z = 0$ mm). A reflecting band-pass filter positioned just before the camera acts as a notch filter in the wavelength range of 500 - 560 nm, excluding any light originating from the 532 nm laser interferometry during the experiment to be recorded.

Optical self-emission images typical for these experiments are presented in Fig. 5.4. The image on the left-hand side shows a background image taken prior to the shot when there was no plasma present. In this case, the targets were illuminated by an LED lamp positioned on the top chamber window. On the right-hand side, a frame captured at 169 ns (431 ns) after the main X-ray pulse is displayed. The X-ray radiation arrives from the top of the images shown in Fig 5.4, illuminating the targets positioned 28 mm away from the X-ray source (central axis of the wire array). This radiation causes material ablation from the surface of the two silicon targets, resulting in the formation of supersonic counter-streaming plasma flows expanding predominantly in the direction normal to the target surface, along the horizontal x-axis, towards the mid-plane of the field of view. Upon collision of these plasma flows, a shocked dense layer is formed at the mid-plane, located between the targets.

As mentioned before, all the data presented in this study are referenced at the implosion time of the array when the main X-ray pulse is generated. This choice is made because it is not expected for any significant plasma material to ablate from the target during the early X-ray emission of the ablating wires or the X-ray precursor. This can indeed be observed in the early-time optical self-emission images captured in the s1203_21 shot, as shown in Fig. 5.5. The initial two frames were captured prior to the complete implosion of the wire array, at a time of 270 ns after current start, which is before the onset of the main x-ray pulse. At this stage, no emission is seen on the target surface indicating a very low level of plasma formation prior to implosion. Instead, only ablated materials from the edges of the targets, as well as CH material from the target holder due to exposure by low X-ray radiation emission, are visible, likely originating from the ablating wires. To observe a significant plasma from the targets, it is necessary to wait until the emission of the main x-ray pulse, as illustrated at 1 ns (250 ns) after the x-ray pulse.

The overall evolution of the plasma flows in a single experiment can be seen in Fig. 5.6, which shows a series of twelve optical self-emission images captured from the end-on view (xy-plane). The

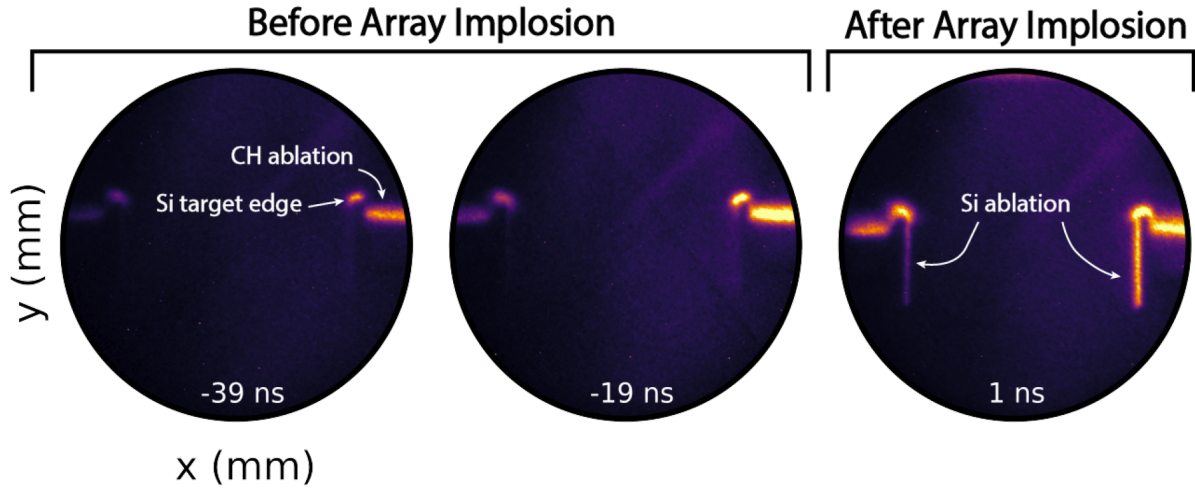


Figure 5.5: Early optical self-emission images from s1203_21 shot taken at three different times. Time equal to 0 ns corresponds to the start of the main X-ray pulse. Prior to the main X-ray pulse, no emission is seen from the target surfaces while plasma appears visible 1 ns after the X-ray pulse.

first frame is recorded 94 ns (356 ns) after the initiation of the main x-ray pulse. The ablating silicon targets are clearly visible at the two sides of each frame, accompanied by the strong emission from the targets corresponding to the counter-streaming plasma flows expanding towards the mid-plane between the targets. This emission lasts for the entire duration of the experiment. At the centre of the images, between the two ablating silicon targets, a strongly emitting interaction region is formed with an initial length (in the y direction) comparable to the lengths of the two targets ($\simeq 10$ mm). We will refer to this structure as "interaction region", "stagnation layer" or "shocked layer", interchangeably. In the early times, the emission profile of the stagnation layer appeared "hollow" with a stronger emission at the boundaries and weaker in the centre, suggesting the stagnation layer is bounded by two reverse shocks. The thickness of this shocked layer (in the direction of the plasma flows) increases with time and the low emission intensity in the middle of the layer becomes less obvious. This "hollow" profile is only observed in the initial four frames of the sequence, up to 169 ns (431 ns) after the x-ray pulse. The layer also increases in length (perpendicular to the flow direction), exceeding the length of the targets. This could be caused by both plasmas escaping the layer driven by the internal layer pressure or due to the existence of diverging plasma flows coming from the targets and interacting along the mid-plane.

As more plasma accumulates in the central region of the layer, the layer width increases, and it expands upstream towards the targets. This layer expansion happens symmetrically relative to the

s0201_22 - Si - Si Colliding Jets - Parallel B-Field - Optical Framing Images

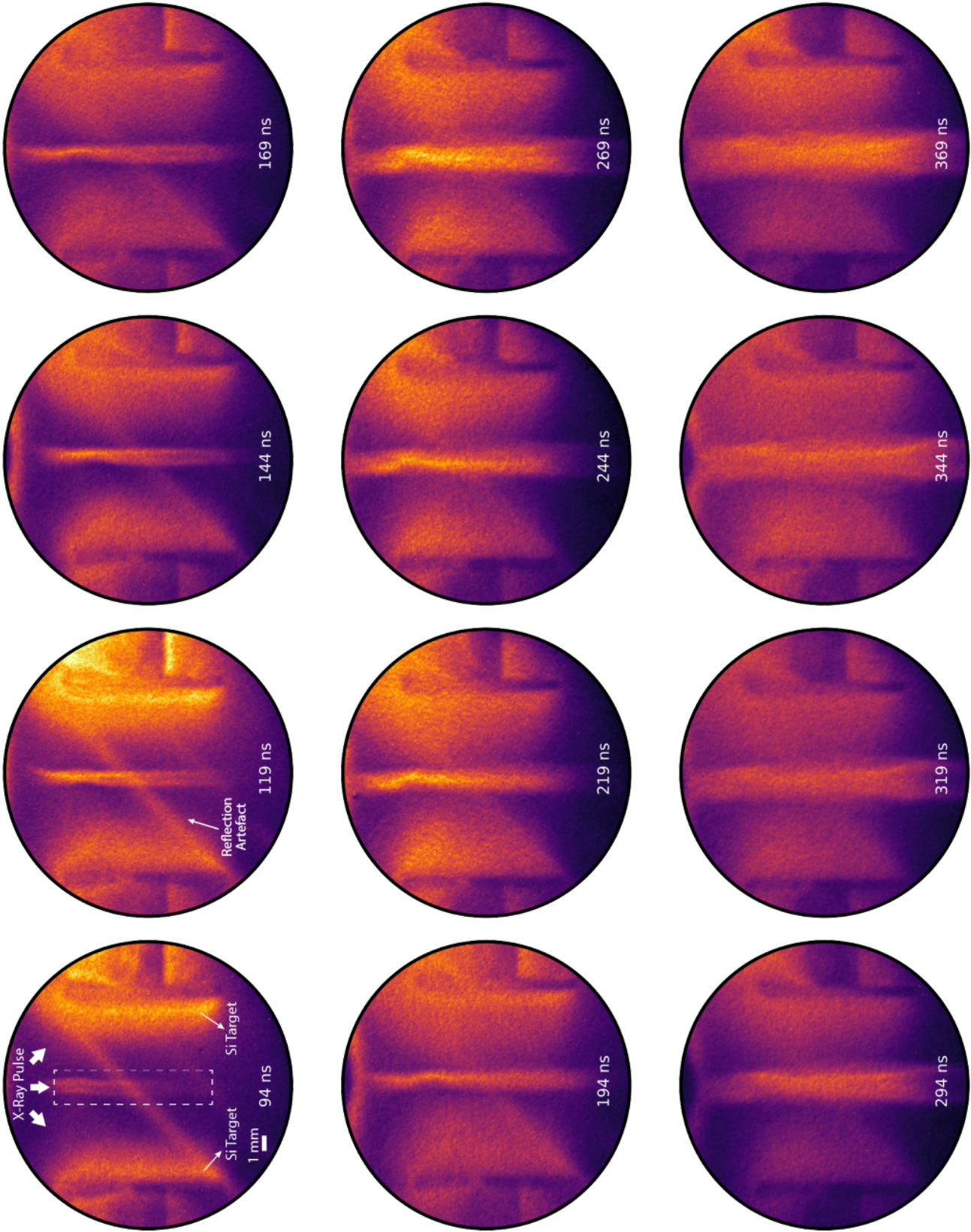


Figure 5.6: Fast framing images taken with 25 ns inter-frame starting from 94 ns (356 ns) after the X-Ray pulse

mid-plane between the targets, suggesting the presence of equal supporting ram pressures generated by the two streaming flows. The entire interaction layer maintains a stable shape until the final captured frame at 369 ns (631 ns), with no distinct internal structures observed in the optical self-emission images throughout the evolution of the system.

We should note the presence of an unfortunate artefact in the form of a light reflection which can be seen in the first few frames, spanning the entire interaction region at a 45-degree angle from left to right. This unwanted reflection persists until the fourth frame and is likely caused by light reflecting off the surface of the end-on sacrificial glass present during the shot.

Despite the limitations of optical self-emission imaging, such as low resolution and uncertainties in the depth of field, valuable quantitative insights can still be gained by analysing how intensity features evolve over time. In particular, the speed of the bounding reverse shocks can be estimated from measurements of the shocked layer width. To determine the layer width from the self-emission data the following steps have been performed. First, a horizontal intensity profile between $x = -5$ mm and $x = 5$ mm is taken by averaging the intensity values over 0.5 mm in the y -direction, around $y = 0$ mm. This line-out, labelled as L0, provides a representative profile of the interaction region, which appears as a distinct peak in the intensity distribution plotted in Fig. 5.7(b). Second, the weighted average value of the distribution is subtracted from the intensity profile such that the peak is aligned with the x -axis origin. The normalised cumulative intensity distribution (red line) is computed from the intensity values between the two minima, which defined the shocked layer boundaries (dark blue line). Finally, An estimate of the shocked layer width, ξ , is obtained by calculating the spatial range between the points at 10% and 90% of the cumulative distribution. This analysis was repeated at different y -positions along the layer resulting in negligible variations in the measured layer width. The results of the measurements of the layer width are illustrated in Fig. 5.7(c), providing insights into the temporal evolution of the shocked structure. It is important to note that the uncertainties associated with these estimates are mainly associated with line-integration effects and possibly the non-optimal position of the camera focus plane. The uncertainties in the measurements are estimated to be approximately 20%, which was quantified by varying the initial percentile range and visually comparing the new values of ξ with the relative position of the layer boundaries in the self-emission images.

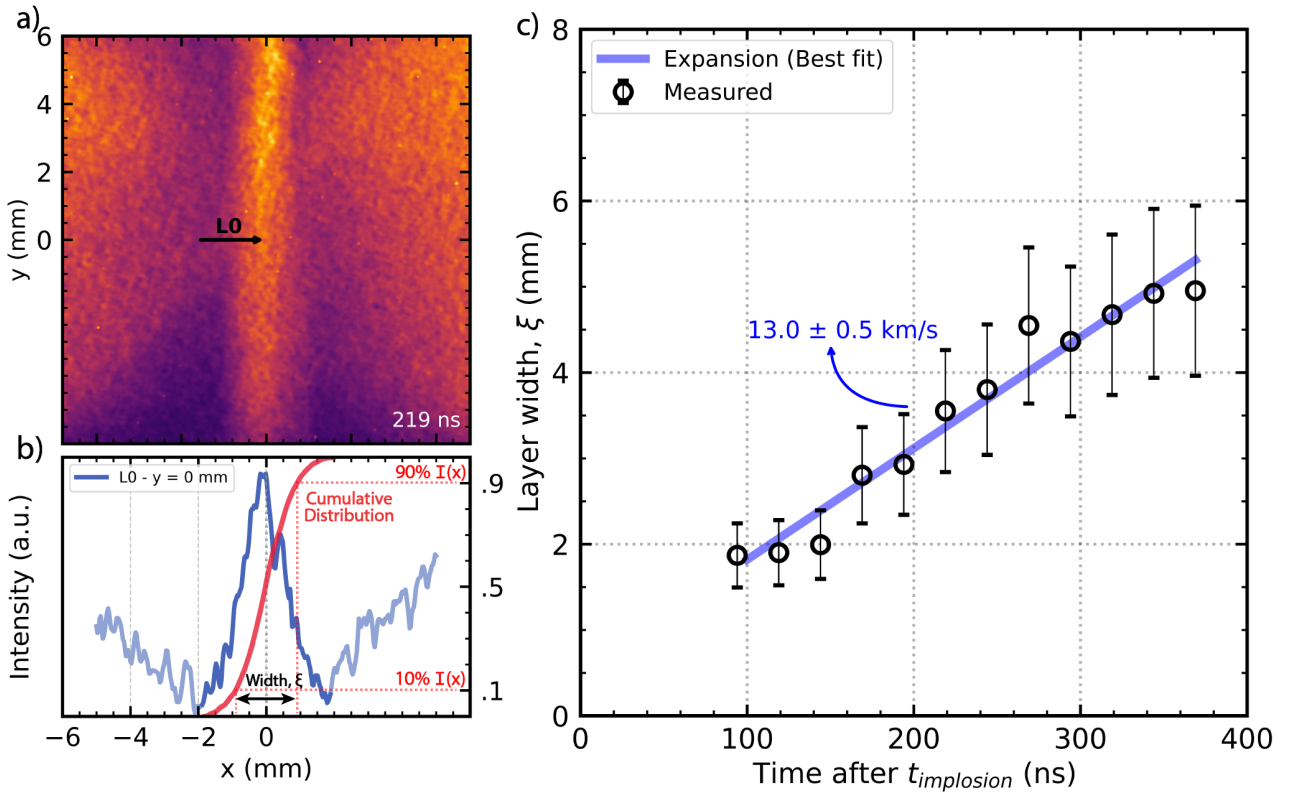


Figure 5.7: (a) Example of optical self-emission frame taken at 219 ns after the x-ray pulse. (b) Intensity profile across the shocked layer from the horizontal line-out (L0) at $y = 0$ mm, extending from $x = -5$ mm to $x = 5$ mm (solid blue line). The cumulative distribution of the intensity profile around the main peak is used to estimate the layer width, ξ by taking the range between the 90 and 10 percentile. (c) Layer width, ξ , measurements as a function of time after X-ray pulse start. The uncertainty was quantified by varying the initial percentile range and visually comparing the new ξ values with the position of the layer boundaries in the self-emission images. The best fit indicates the average expansion velocity of the layer.

As can be seen from the plot shown in Fig. 5.7(c), the interaction layer expands with time as more shocked material accumulates and the bounding shocks propagate back into the upstream plasma flows. During the early stages of the interaction region formation (the first $\simeq 50$ ns), the layer width remains approximately constant at $\xi \simeq 2$ mm, and then expands at an approximately constant speed, eventually reaching its maximum width of $\xi \simeq 5$ mm in the last frame at 369 ns. This corresponds to a layer expansion speed of $V_{layer} = 13 \pm 0.5$ km/s. The velocity of the shocks in the lab frame can be estimated by taking half the speed of the layer expansion, given by $V_s \simeq \Delta\xi / 2\Delta t$, where $\Delta\xi$ represents the change in layer width over the corresponding change in time, Δt , giving reverse shock velocity in the lab frame, $V_s \simeq 7$ km/s.

In summary, optical self-emission imaging allowed to study the dynamic evolution of the interaction

region resulting from the collision of two counter-streaming plasma jets. A sequence of 12 images was captured in a single shot, with a 5 ns exposure time and an inter-frame separation of 25 ns. The first image was recorded 94 ns (356 ns) after the implosion of the array, already revealing the presence of the interaction region. This is reproducible for images taken from different shots, where the layer becomes detectable between $\simeq 100$ ns and $\simeq 120$ ns after the X-ray pulse, indicating an average flow velocity of $V_{flow} \simeq 60$ km/s. The interaction region is bounded by features which can be interpreted as two planar reverse shock structures that are uniform along the y-axis. Initially, the layer had a width of approximately 1 mm, and a narrow slab of low plasma emission was observed between the two bounding shocks. As time progressed, this low-emission slab disappeared, leading to a shocked layer with uniform emission. Notably, the expansion of the shocked layer proceeded at an approximately constant rate, $\simeq 200$ ns from the start of the X-ray pulse. The expansion phase exhibited symmetrical behaviour, with both reverse shocks propagating at an equal velocity of $V_s = 7$ km/s in the laboratory frame. Further discussion of these observations will be presented later in the chapter.

5.1.2 Electron density measurements

Accurate measurements of the electron density play a crucial role in the investigation of shock systems, as they facilitate the characterization of upstream and downstream shock conditions (such as the compression ratio). Coupled with velocity measurements, these density measurements enable the calculation of the ram pressure exerted by the system. In the present experiment, a 532 nm laser imaging interferometry setup is utilized to image the plasma system. Specifically, the laser beam is directed into the plasma along the z-direction, while the imaging takes place in the xy-plane, providing an end-on view of the plasma flow interaction.

Fig. 5.8 shows the end-on raw background and shot interferograms obtained using a 532 nm probing laser. The shot interferogram is captured at 218 ns after the X-ray pulse, which corresponds to 480 ns after the start of the current. The interferograms are oriented such that the X-ray radiation propagates from top to bottom, with the Z-pinch located above the top of the images. The origins of the x and y axes are centred at the mid-plane between the targets and at half-target length, respectively. The two targets are positioned at a measured distance of 13 mm from each other and $\simeq 30$ mm from

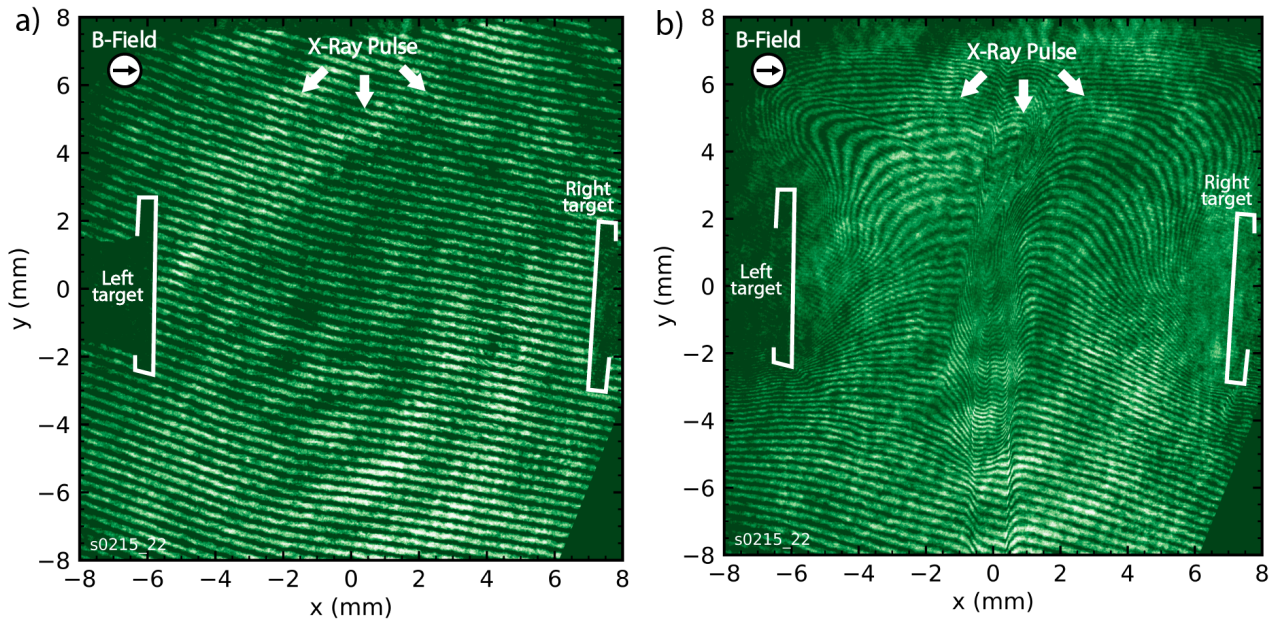


Figure 5.8: End-on 532 nm raw interferograms from the shot s0215_22. The X-ray source is located above the images, outside of the camera field of view. (a) Background interferogram taken before the experiment, in the absence of plasma. (b) Shot interferogram at 218 ns after the main X-ray pulse. Displacement of fringes is the strongest at both target surfaces and also in the central region at the position of the shocked layer.

the wire array. The background interferogram, as shown in Fig. 5.8(a), exhibits evenly spaced straight fringes that are nearly perpendicular to the target surface. These fringes have a small angle relative to the flow direction. The chosen orientation of the fringes in the background interferogram maximizes the diagnostic sensitivity to density variations along the x -axis, where important features are likely to appear. Additionally, aligning the fringes parallel to the expected direction of the density gradient on both sides of a transition (e.g. for a planar shock, the fringes should be parallel to the shock normal) makes it easier to track the fringes across a density jump. In this configuration, the fringes should be parallel to the shock normal. In the raw interferogram reported in Fig. 5.8(b), the displacement of the fringes due to the presence of plasma is observed. Fringes can be traced across the entire interaction region and most regions around the targets. Regions of lost fringes near the target surfaces (at $x = -5$ mm and $x = 6$ mm) and in the upper region of the shocked layer, $0 \text{ mm} < y < 5 \text{ mm}$, are either caused by the probing laser being deflected beyond the maximum acceptance angle of the optical system or being absorbed due to inverse bremsstrahlung. These areas are not included in the analysis and are masked out in the final electron density map. Following the same interferogram analysis as presented in previous chapters, the electron density map can be reconstructed by relating the fringe shift of the

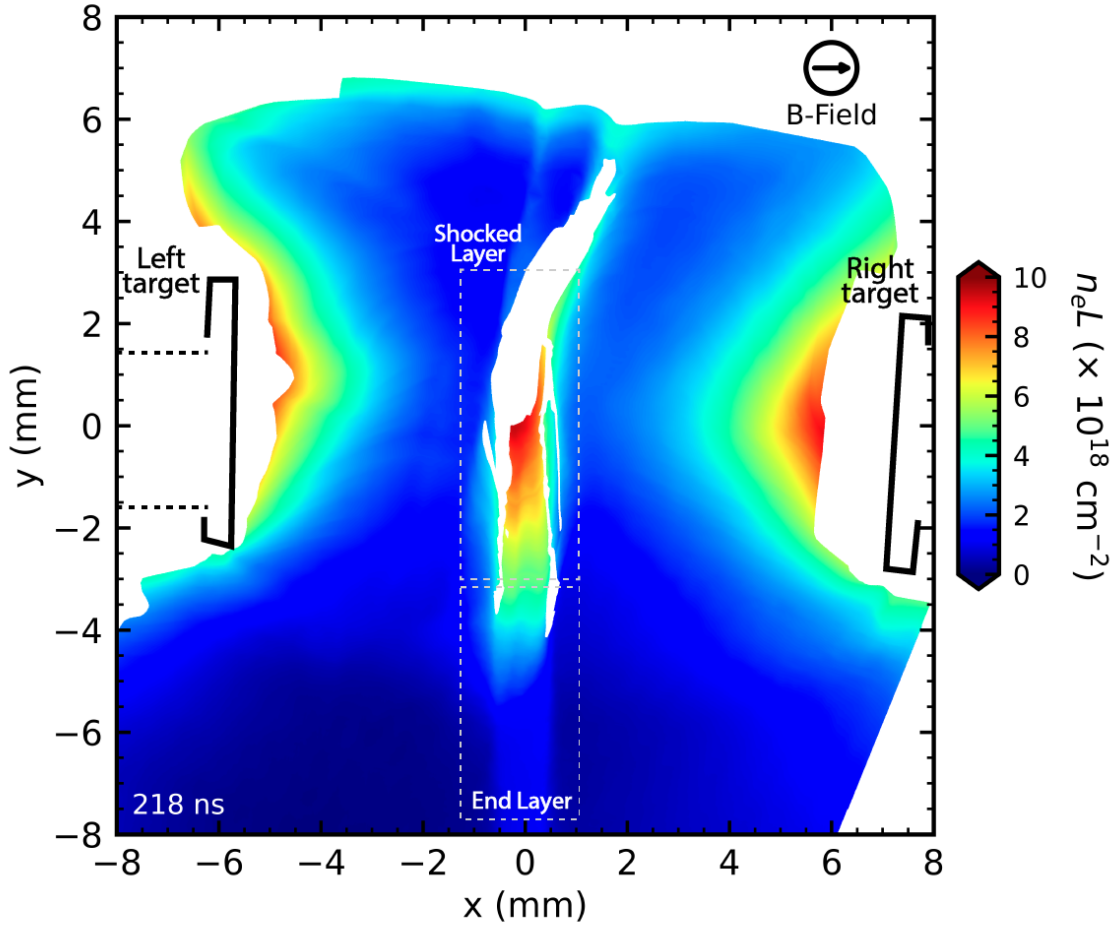


Figure 5.9: Electron density map reconstructed from the raw interferograms presented in Fig. 5.8 and taken at 218 ns (480 ns) after the X-ray pulse. The two targets are represented by solid black lines located at the two opposite extremes of the figure, $x = -6$ mm and $x = 7$ mm, respectively.

interference pattern to the background image. To determine the absolute fringe shift relative to the background interferogram, a region with zero fringe shift must be identified for the shot interferogram. This corresponds to a region in the shot interferogram that does not contain any plasma, which in this case, can be seen in the bottom part of the shot interferogram, $(x, y) = (-4 \text{ mm}, -6 \text{ mm})$ and $(x, y) = (4 \text{ mm}, 6 \text{ mm})$.

The result of the tracing analysis is presented in Fig. 5.9 as the line-integrated electron density, $n_e L$, while the corresponding volumetric density can be estimated assuming that the plasma length in the probe direction is equal to the target length, $L = 5$ mm. It is seen from Fig. 5.9 that the highest electron density is recorded near the surface of the targets. The density then gradually decreases as the plasma flows expand towards the centre, ranging from $n_e \simeq 1.6 \times 10^{19} \text{ cm}^{-3}$ at $x \simeq 4$ mm to

$n_e \simeq 4 \times 10^{18} \text{ cm}^{-3}$ just before reaching the shocked layer. Both streaming flows exhibit a uniform density profile in the y -direction at the upstream flow, just outside of the layer. In the shocked layer which has $\Delta x \simeq 2 \text{ mm}$, the electron density is larger, reaching at the centre a maximum density of $n_e \simeq 2 \times 10^{19} \text{ cm}^{-3}$. The width of the layer measured by the interferogram correlates very well with the thickness of the layer estimated from optical self-emission, $\xi \simeq 1.8 \text{ mm}$, at the time when the interferometry was recorded.

At the very top part of the field of view, $y > 5 \text{ mm}$, this is a region of a relatively low density, not related to the interaction of the flows. This is likely due to plasma material originating from the sacrificial protective glass positioned at the top of the hollow support structure that enters the field of view of the diagnostic. This low plasma material does not affect the system at any point during the experiment as indicated by both the self-emission images and interferometry data, since it is not in the same vertical position where the plasma flows and the shocked layer are present.

In the layer, the largest electron density is concentrated at y -positions aligned with the target extent ($-2.5 \text{ mm} < y < 2.5 \text{ mm}$). For the bottom part of the layer, when measurements are available, it is seen that the density rapidly decreases for $y < -2.5 \text{ mm}$. To further characterise the plasma distribution within the layer, horizontal density profiles, $n_e(x)$, for various y -positions are discussed next.

Fig. 5.10(b) shows three horizontal line-outs taken at different y -positions across the shocked layer, averaged over $80 \text{ } \mu\text{m}$ in the y -direction. The densest measured region within the layer is localised at the axis origins aligned with the targets and extending in y -directions over a length comparable to the target size. This observation suggests that the majority of the stagnant plasma remains confined to the central region of the layer.

For $y > 0 \text{ mm}$, the density of the layer cannot be measured due to stronger density gradients. While direct evidence is currently lacking, a possible explanation is that the presence of higher magnetic pressure on that side of the layer (which is closer to the z -pinch axis) may impede the escape of plasma, resulting in somewhat uneven density distributions within the stagnant layer and the bending of the top part of the layer in the direction along the field lines. An additional reason for the higher density in the top half of the layer is that the targets are more exposed to a slightly stronger X-ray radiation

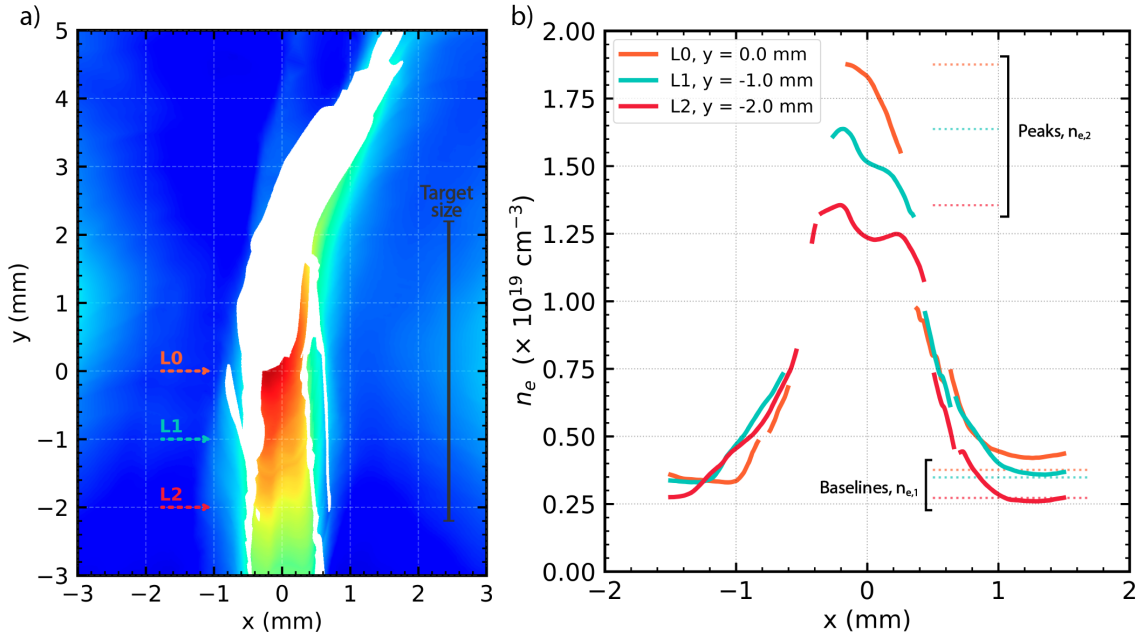


Figure 5.10: Zoomed-in region of the end-on electron density map, $n_e L$, presented in Fig. 5.9. (a) View of the shocked layer region co-aligned with the silicon targets. The densest region is confined within an area of the length size of the target, $\simeq 5$ mm. (b) Three density profiles, $n_e(x)$, taken from four horizontal line-outs at different y -positions.

flux than the bottom part. The amount of X-ray flux reaching the target is approximately proportional to the inverse of the distance, $\propto 1/d$, (in cylindrical approximation) since the wire array is much taller in the vertical direction ($\simeq 22$ mm) than the target height (5 mm). The difference in X-ray flux between the top edge ($y = 2.5$ mm) and the bottom edge of the target ($y = -2.5$ mm) can be estimated at approximately 20%. Although relatively small, this can potentially contribute to the appearance of large density gradients in the top part of the shocked layer.

In the region of the layer where density measurements are fully resolved ($y < 0$ mm), it is seen that the electron density decreases along the layer (along y), perpendicularly to the plasma flow direction. This decrease is relatively slow in the region corresponding to the target size as can be seen from the decrease of the maximum density profiles shown in Fig. 5.10(b). After that, there is a faster and stronger reduction of the density. Two possible scenarios can potentially explain the observed mass distribution in the layer: the shocked plasma moves away from the central stagnant region in the direction along the layer due to the gradient of internal pressure and/or due to the presence of diverging plasma flows coming from the targets at $-8 \text{ mm} < y < -4 \text{ mm}$.

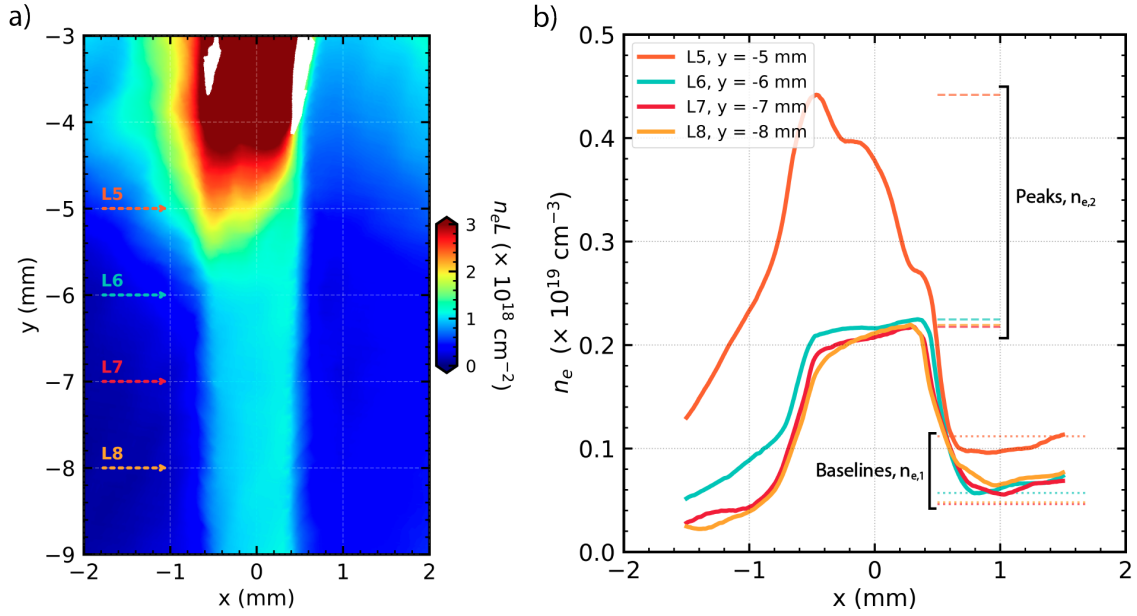


Figure 5.11: (a) Zoomed-in and rescaled heat-map of the shocked layer between $-3 \text{ mm} < y < -9 \text{ mm}$. (b) Four horizontal density profiles taken at different y positions along the layer.

The zoomed-in and re-scaled heat-map (from 0 to $3 \times 10^{18} \text{ cm}^{-2}$) of the bottom part of the density map shown in Fig. 5.10 is shown in Fig. 5.11, revealing more details of the density distribution. Horizontal line-outs taken at various y positions show the density variation along the layer, in the region between $-9 \text{ mm} < y < -3 \text{ mm}$, Fig. 5.11(b). It is seen that there is a sharp density decrease, by a factor of 2, between the profile at $y = -4 \text{ mm}$ and the profiles at $y = -5 \text{ mm}$, $y = -6 \text{ mm}$, $y = -7 \text{ mm}$. It is seen that density profiles do not change within the layer for $y \leq -5 \text{ mm}$. These observations can be interpreted as the presence of both reasons mentioned above for the layer expansion along the y -direction. For $y > -5 \text{ mm}$ there could be a significant (dominant) contribution from the thermal pressure gradient along the layer. For $y \leq -5 \text{ mm}$, it is most probable that this region is formed by stagnation of the diverging, not fully confined by the magnetic field, plasma flows from the targets.

To investigate earlier stages of the stagnation layer formation and to compare the density profile with the self-emission profiles recorded with the optical fast frame camera (Fig. 5.6), we performed interferometry measurements at 174 ns after the start of the X-ray pulse. The raw interferogram and the reconstructed electron density map are shown in Fig. 5.12. The origins are positioned at mid-distance and at the centre of the targets, such that the interferograms show only the bottom half of the shocked layer. The interaction layer extends along the y -direction with a length comparable to the size of the

two targets, $\simeq 5$ mm, but only half of the stagnation layer is shown.

The interferogram and the obtained density map show a number of very interesting features, most clearly those can be seen on horizontal density line-outs presented in Fig. 5.12(d). Firstly, the density profiles show density reduction in the middle which is consistent with the "hollow" emission profile seen in the optical self-emission images. The recorded density reduction at the centre of the layer is relatively small, at best 8% for the profile at $y = -0.5$ mm. Second, these line-integrated measurements allow to obtain an upper limit on the shock width. The total layer width is approximately $\simeq 1$ mm with the shock transitions happening at a spatial scale of $\simeq 0.6$ mm. The density distribution, $n_e(y)$, also changes along the shocked layer. The maximum electron density is $4.5 \times 10^{18} \text{ cm}^{-3}$ at $y = -1$ mm, dropping to $3 \times 10^{18} \text{ cm}^{-3}$ at $y = -2.5$ mm and continuously reducing in the region outside of the target size. Third, a "precursor" density front can be observed in the upstream flow (see Fig. 5.22), extending by $\simeq 0.4$ mm and visible in both raw interferograms and density profiles. This density precursor will be discussed later in the chapter.

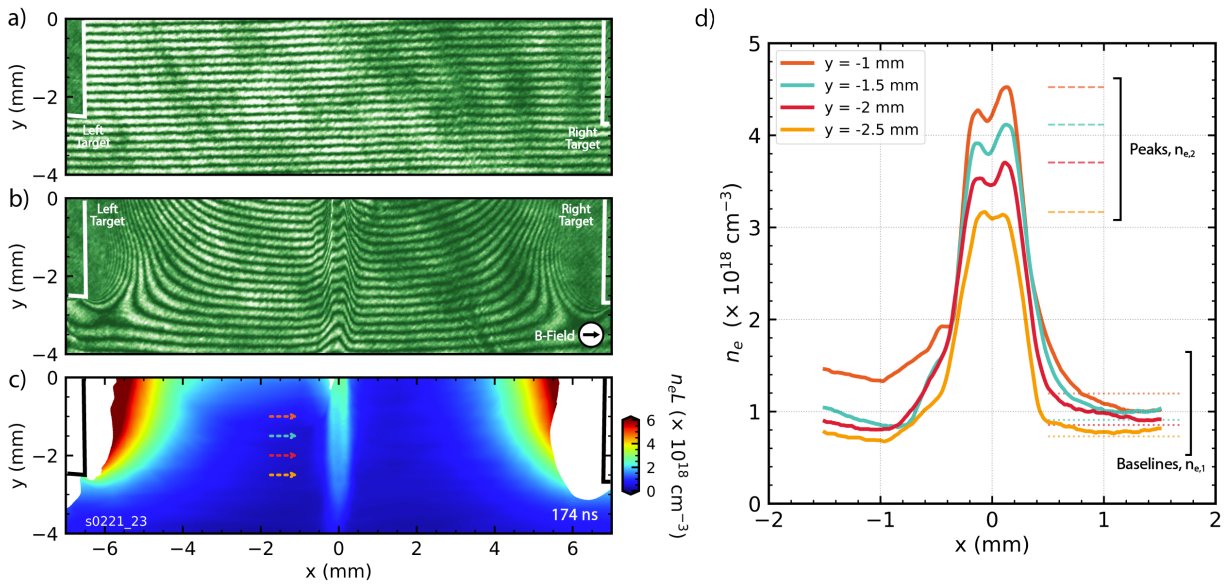


Figure 5.12: (a), (b) Background and raw interferograms from an early time shot (s0221_23) at 174 ns after the X-ray pulse. (c) Reconstructed line-integrated electron density map. (d) Four horizontal density profiles taken at different y positions along the shocked layer.

The measured density profile allows to determine the electron density jump of the reverse shock. The compression ratio is defined as the ratio of the post-shock electron density to the pre-shock density, $R_{n_e} = n_{e,2}/n_{e,1}$, where the subscripts 1 and 2, represent regions outside and inside of the shocked layer,

respectively. This is not a mass density compression ratio but can be considered as an upper limit of it. Based on the electron density profiles presented in Fig. 5.10(b), the determination of the compression ratio, R_{n_e} , strongly relies on its definition. In this case, an effective approach is to calculate the ratio between the peaks and the corresponding baselines for each electron density profile. This provides a range of possible values for R_{n_e} . The analysis reveals an average compression ratio of $R_{n_e} = 4.5 \pm 0.5$ taken from the region co-aligned with the targets, while, from the density profiles at $y \leq -2.5$ mm, the measured density jump is lower, $R_{n_e} = 3.5 \pm 0.5$. The implications of these observed electron density jumps will be discussed later in the chapter.

5.1.3 Thomson scattering measurements of velocity and temperature

In this section, the results from the Thomson scattering diagnostic are discussed, providing measurements of the in-plane velocity and temperatures across the shocked layer. The Thomson scattering geometry used in this experiment is introduced as well as the orientation of the scattering vectors. Finally, raw spectrograms and spatially resolved plasma parameters are presented, followed by a discussion of the fitting procedure used to extract velocity and temperatures.

Thomson scattering experimental setup

The Thomson scattering technique has been employed to characterise the shocked layer formed from the interaction of two colliding flows. Similarly to the experiments presented in Chapter 4, the diagnostic is employed in the collective regime focusing on the observation of the ion acoustic feature. In this regime, the scattered spectra are influenced by the electron density, plasma flow velocity, ion temperature, and the product of the ion charge Z and the electron temperature T_e . For more details refer to Chapter 2.

The Thomson scattering experimental setup specifically designed to investigate the interaction region of two colliding plasma flows, is illustrated in Fig. 5.13. The laser beam crosses the shocked layer at an angle of 45° . The position of the laser spot is chosen to be at the mid-plane of the field of view, located 7.5 mm away from the target surfaces, and positioned on the xy plane at half the height of

the targets. The top portion of the figure shows the position of the imploding wire array. The focused laser beam propagates along the direction indicated by wavevector k_{in} . The scattered light is collected from two opposite locations at a 90° angle relative to the probing beam. The collection optics for the two fibre bundles, designated as bundles A and B, consisting of lenses with a focal length of 20 cm positioned at 50 cm from the object and set to image the scattering volumes at a magnification of $M = 0.9$. The 22 collection volumes for each fibre bundle (producing 2×22 TS spectra in total) are spaced $\simeq 200 \mu\text{m}$ apart from centre to centre of the scatter volumes and span a range of 5 mm in the x-direction. This arrangement allows for the structure of the stagnation layer to be resolved, while also collecting multiple measurements of the pre-shock (upstream) plasma flows. Fibres from each bundle record Thomson scattering light from the same locations along the focused laser beam. This advantageous configuration ensures that the spectrum of the scattered light emitted from each plasma volume is collected twice, once from fibres in bundle A and again from the corresponding fibres in bundle B.

The described Thomson scattering geometry is characterised by the vector diagram illustrated in Fig. 5.13(Right). The collection directions are represented by the vectors $\mathbf{k}_{out,A}$ and $\mathbf{k}_{out,B}$, which are orthogonal to the incident laser beam $\mathbf{k}_{in}$. This geometry allows measuring plasma flow velocities along the two scattering vectors, $\mathbf{k}_A$ and $\mathbf{k}_B$ from the Doppler shift of the Thomson scattering spectra. Vector $\mathbf{k}_A$ is parallel to the streaming plasma flow (parallel to the shock normal) while vector $\mathbf{k}_B$ is along the interaction region (shown in grey). Each fibre bundle will, therefore, measure the corresponding components of the velocity vector in the xy-plane from the same positions. The fitting of the spectra also provides information on the plasma temperatures.

Experimental results

The Thomson scattering measurements were obtained in the same experiment and at essentially the same time ($t_{TS} = 221(483 \text{ ns}) \text{ ns}$, $t_{interf} = 218(480 \text{ ns}) \text{ ns}$) as the interferometry measurements of electron density distribution presented in Fig. 5.9. The exact positions of the Thomson scattering volume with respect to the density map were determined using the technique described in Chapter 2. These Thomson scattering volumes are overlaid on the density map as shown in Fig. 5.14.

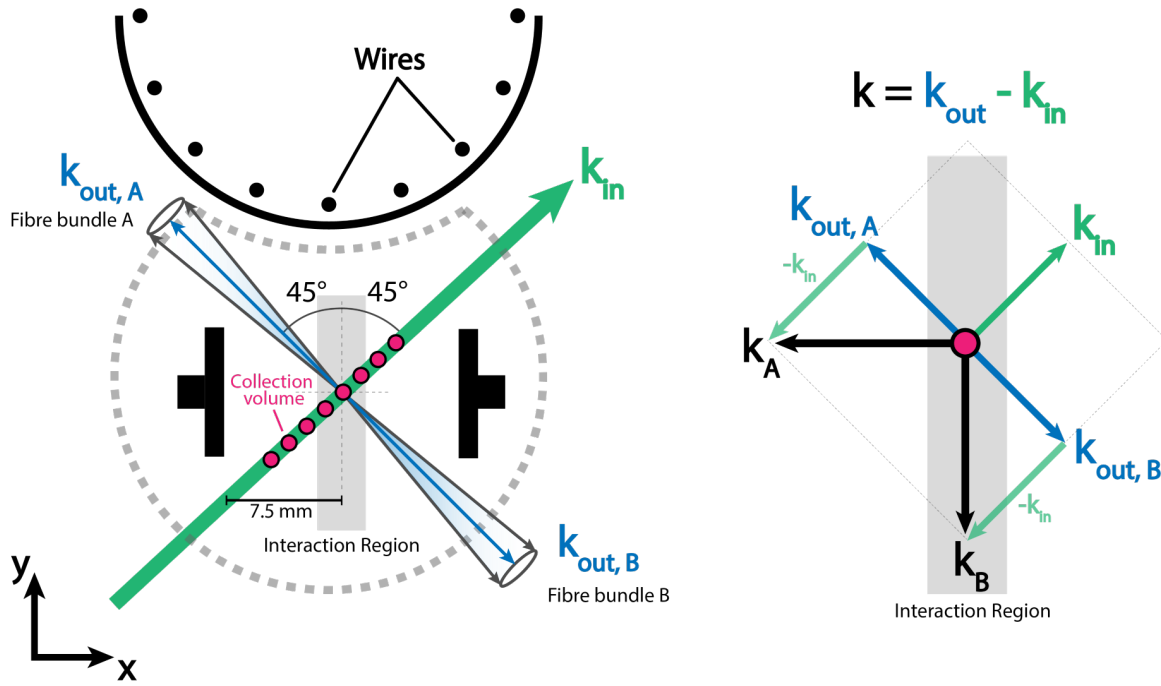


Figure 5.13: (Left) Schematic illustration of the Thomson scattering setup (not to scale). The 532 nm laser beam is focused and propagates through the interaction region (grey) from the bottom to the top of the figure, forming a 45° angle with the central axis of the shocked layer. Seven out of the 44 scattering volumes along the path of the laser are shown. The scattered light is collected from two fibre bundles, bundle A and bundle B, positioned in the x-y plane at a 90° angle relative to the laser beam. The separation between each scattering volume is approximately $200\ \mu\text{m}$. (Right) Vector diagram of the Thomson scattering geometry. The orientation of scattering vectors, k_A and k_B , allow measurements of the plasma flow velocity vector in the xy-plane.

The measurements cover a spatial range of approximately $\simeq 5\ \text{mm}$ in both the x and y directions, crossing the interaction region diagonally near the origins of the axes. The scattering volumes are numerically labelled from 1 to 22, with 5 collecting volumes located on the bottom-left side of the shocked layer, 5 volumes within the stagnation region, and the remaining 12 volumes on the upper-right side of the layer. Fitting the Thomson scattering spectra collected from various locations allows to measure (or determine) T_i and $\bar{Z}T_e$. This fitting process is constrained by using the information on the electron plasma density at each corresponding location obtained from the interferometry measurements.

The raw spectrograms obtained from fibre bundles A and B are presented in Fig. 5.15, accompanied by two examples of Thomson scattering spectra. It can be seen in both spectrograms that the Thomson scattering signals are concentrated within a narrow range of approximately $3\ \text{\AA}$ near the 532 nm laser wavelength, exhibiting the characteristic double peak of ion-acoustic waves and good overlap in fibre positions. Overall, spectral signals are detected in most of the collecting scattering volumes outside

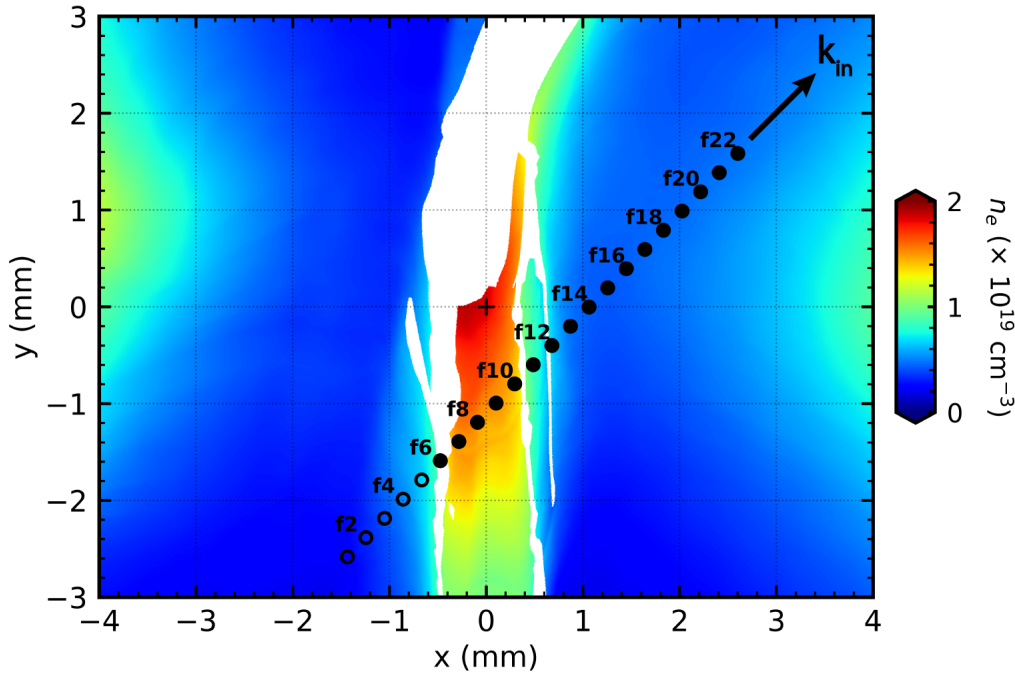


Figure 5.14: End-on electron density map at 218 ns (480 ns) with the scattering volumes positions. The accurate location of each scattering volume is estimated by interpolating the positions of the scattered light of a metallic object along the laser path in pre-shot images.

and inside the shocked layer.

Examining the spectrogram of bundle A, it is seen that for fibre 11 ÷ 22 spectra are significantly Doppler shifted. These fibres correspond to positions outside of the stagnation layer and the recorded blue shift of the Thomson scattering signals indicates the direction of the flow, from the target towards the shocked layer, as expected. Fibres from 10A to 7A show very small Doppler shifts indicating low flow velocity in the layer. The transition from high to low velocity is very localised, happening between fibre 10A and 11A, on a scale smaller than the scattering volume separation, $\approx 200 \mu\text{m}$. Deeper into the shocked layer, the signals weaken, with spectra from fibres 7A and 8A being just above the self-emission background. This, together with the loss of Thomson scattering signals from fibre 6A to fibre 1A, could be possibly due to a misalignment between the laser beam and the collecting fibre positions for this region.

The majority of spectral signals from bundle B are captured from plasma volumes between fibre 1B ÷ 12B. A small Doppler blue shift is recorded for these spectra, which remain almost constant for the entire spatial range, outside and inside of the shocked layer. For this observation direction,

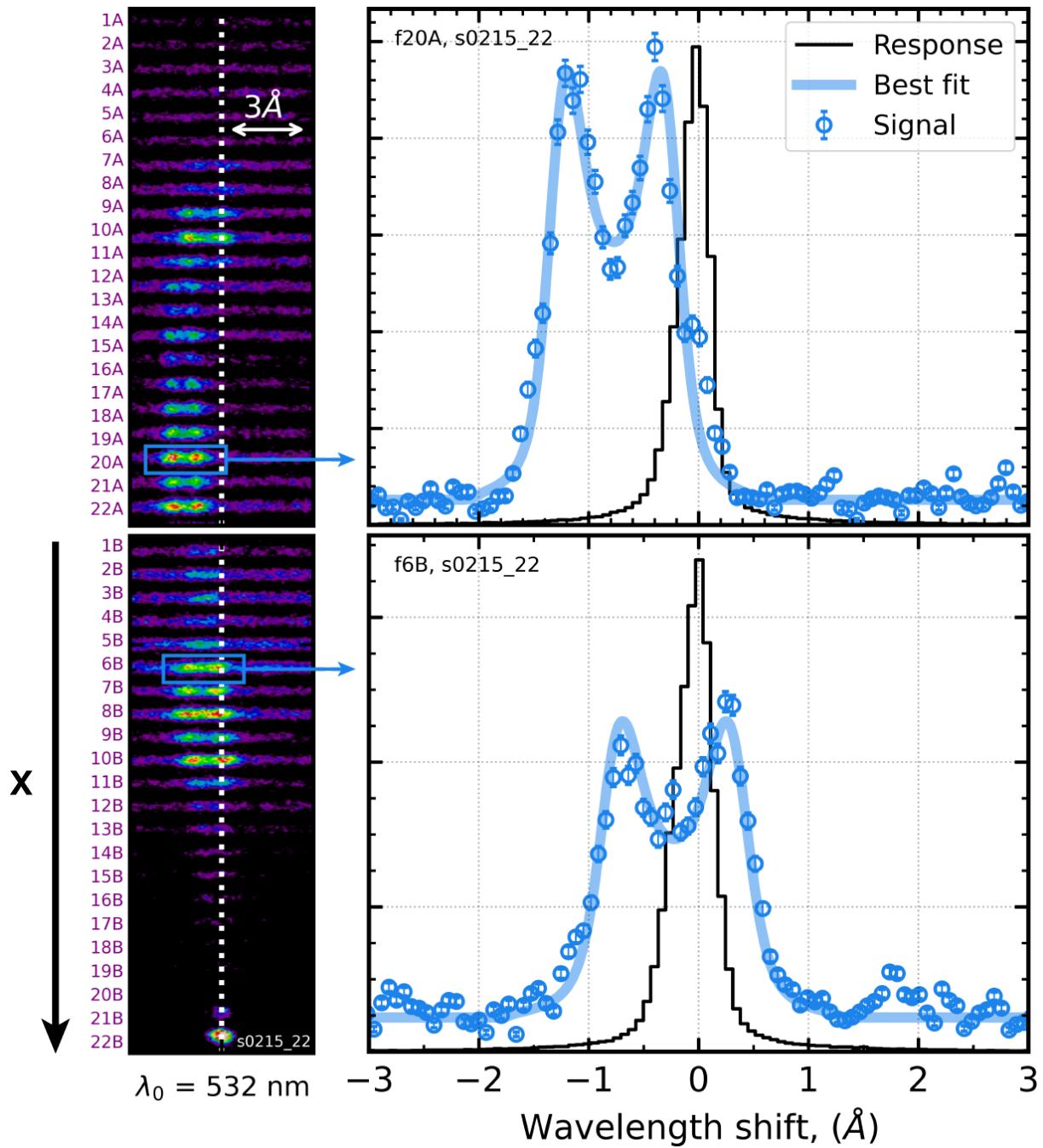


Figure 5.15: (Left) Thomson scattering raw spectrograms from s0215_22 taken at 221 ns (483 ns) after X-ray pulses. The top panel represents fibre bundle A, while the bottom panel represents bundle B. The laser wavelength is indicated by the dotted white line. Most of the recorded spectra display ion acoustic features. Scattering signals are absent from fibre 1A to 6A along the $\mathbf{k}_A$ vector and from 13B to 22B along the $\mathbf{k}_B$ vector. (Right) Examples of spectra from fibre f20A (top) and f6B (bottom). The top spectrum corresponds to a measurement taken outside the shocked layer, while the bottom spectrum represents a measurement taken inside the shocked layer. The black line represents the response of the Thomson scattering diagnostic in the absence of plasma.

the measurements are only sensitive to velocity along the stagnation layer, along the $\mathbf{k}_B$ vector. For fibres within the shocked layer, $6B \div 10B$, this suggests non-zero velocity along the layer. For fibre numbers greater than 12, outside the layer (Fig. 5.14, $x > 0.5$ mm), there is almost no signal detected from bundle B, indicating either laser probe deflections due to density gradients or misalignment of the beam with the collecting fibres. A quantitative discussion of the plasma flow velocity will be presented later in this section.

In addition to the Doppler shifts, the broadening of Thomson scattering spectra is observed for spectra within the stagnation layer, for both bundles, A and B. The fitting of the Thomson scattering form-factor $S(k, \omega)$ (see Chapter 2), to determine plasma temperatures was performed for all the recorded spectra, and two examples from fibre 20A and 6B are shown in Fig. 5.15. The Thomson scattering signals are represented by the shaded blue line, derived by summing up over camera pixels of one fibre in the vertical direction of the spectrogram, to boost the signal-to-noise ratio against self-emission. The best fit of the theoretical spectrum to the data is shown as a solid blue line, while the solid black line shows the response of the spectrometer. Since the width of the spectral response function is non-negligible, $2\sigma_\lambda \simeq 0.5 \text{ \AA}$, the theoretical fitting includes the convolution with the spectrometer response. In both examples of spectra (from fibres 20A and 6B) the Doppler shift is clearly visible, allowing for accurate determination of plasma velocity. In addition, the spectra are characterised by two well-separated ion acoustic peaks from which the product $\bar{Z}T_e$ can be determined. The similarity in width between the peaks and the spectrometer response, however, poses a challenge for precise measurement of the ion temperature (T_i) which will be discussed later in the section. The spectral intensity of plasma self-emission is relatively small and it extends over a range of wavelengths much broader than the width of the ion acoustic spectra. This background signal appears noisy but uniform and can be averaged and subtracted from the measurements.

Velocity measurements

The results of the Thomson scattering analysis, allowing measurements of the plasma flow velocity from the Doppler shifts, are summarized in Fig. 5.16. The spatially resolved velocity components, V_x and V_y are plotted in Fig. 5.16(a) and (b) for bundles A and B, respectively. The shaded grey region represents the position of the shocked layer. The error bars associated with the velocity measurements

are small (± 1 km/s) and fall within the marker size, indicating a high accuracy of these measurements. Outside the shocked layer, V_x slightly increases along the flow from 30 km/s at $x = 2.6$ mm to 35 km/s at $x = 0.8$ mm closer to the shocked region. Velocity measurements along k_A , outside of the stagnation layer, are only available on the right-hand side of the layer, specifically from fibre 22A to 11A, as there were no recorded spectra from fibre 1A to 6A. Although direct measurements of V_x from the plasma flow generated by the left silicon target are lacking, the high symmetry of the system, as indicated by the optical self-emission images (Fig. 5.6), suggests that similar values of V_x observed from fibre 11A to 15A can be expected in this case. At the shock transition, there is a sharp and significant reduction in V_x from 30 km/s to 5 km/s over one collection volume (between fibres 11A and 10A). Within the shocked layer (from fibre 10A to fibre 7A), V_x remains on the whole, non-zero, with an average speed of approximately 5 km/s. This motion of the plasma is difficult to attribute to the rigid displacement of the layer, given the symmetric expansion of the layer observed from optical self-emission images (Fig. 5.6).

At this stage, it is interesting to compare the density profile measured by the interferometry with the velocity V_x . It can be seen that the velocity V_x follows an inverse profile in comparison to the electron density profile obtained from a line-out along the scattering volumes with a sharp velocity transition at the boundary of the shocked layer, occurring over a single inter-fibre distance of $\simeq 200$ μm instead of the observed density jump at a larger spatial scale, $\simeq 0.6$ mm. It is important to note that this velocity transition most likely happens at much shorter scales due to high plasma collisionality which will be discussed later in the chapter.

Velocity measurements along k_y from bundle B are presented in Fig. 5.16(b). The range of available measurements spans from $x = -1.4$ mm to $x = 1.2$ mm, covering regions within and outside the layer on both layer sides. It is seen that the measured V_y component is small, approximately -5 km/s, and similar for all the Thomson scattering positions. This means that the plasma inflow is predominantly in the direction of the target normal to the layer. Some divergence of the flow can be seen in regions outside of the target extent, -2 mm $< y < -3$ mm, with a maximum diverging angle of $\simeq 9^\circ$. There is no noticeable increase of V_y in the layer at the position of the scattering volumes in comparison to the values of V_y outside the layer. This suggests that there is no significant acceleration of plasma in the direction along the layer and in the central part of the layer, for up to 1.5 mm in the y -direction

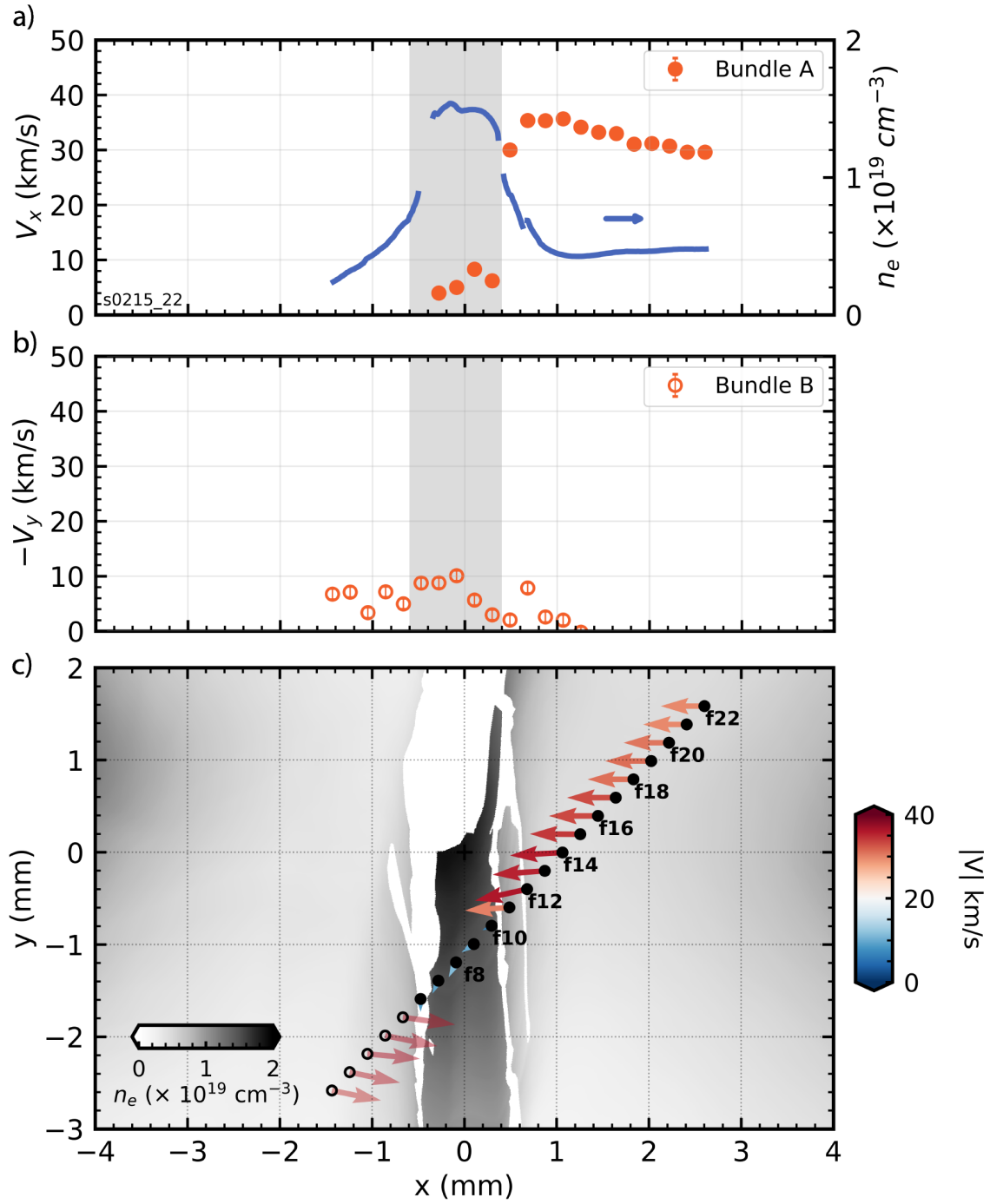


Figure 5.16: Thomson scattering results for the s0215_22 dataset obtained at 221 ns (483 ns) after the start of the X-ray pulse. (a) and (b) depict the velocity components, V_x and V_y , respectively, where V_x is parallel to the x-axis and V_y is perpendicular to it. The error in determining the Doppler shift is negligible, with the error bars being well within the marker size. (c) presents the comparison between the electron density profile and the reconstructed in-plane velocity profile. (d) a vector representation of the in-plane velocity is shown. For fibres 1 to 5, the assumption is made that V_x is equal to the values measured from fibres 11 to 15.

from the middle.

The above observations are further illustrated in Fig. 5.16(c) by the representing vectors of $\mathbf{V}$ by combining the two components V_x and V_y . The velocity vectors for each fibre are indicated by arrows, with their length proportional to the calculated speed, along with the corresponding heat-map of the velocity module. As discussed earlier, velocity vectors in the plasma inflow predominantly point along the x-direction, towards the central region. The shaded velocity vectors from fibres 1 to 6 are estimated by assuming V_x to be equal and opposite to the measured values at fibres 11 to 16. Within the shocked layer, the magnitude of the in-plane velocity significantly decreases, and its direction is primarily redirected along the shocked layer.

Temperature measurements

In addition to inferring plasma velocity from Doppler shift measurements, the characteristic plasma thermal properties, ZT_e and T_i are determined by fitting the theoretical form factor $S(\mathbf{k}, \omega)$ of the scattered spectrum, as described in Chapter 2. The recorded spectra exhibit well-separated ion acoustic peaks, indicating that the ionization-electron temperature product, ZT_e , is much larger than the ion temperature, T_i . This means that T_i is expected to be small and ZT_e can be estimated with good accuracy.

To separate Z and T_e from the measured ZT_e product, an atomic model (FLYCHK atomic code) is used. While these values can be determined with reasonable accuracy, the estimate of T_i remains challenging if used as an independent parameter. This is mainly due to the fact that the broadening of the peaks is comparable to the diagnostic response function. The associated error bars reported in the temperature measurements represent an estimation made during the fitting process, involving manual parameter variation and visual inspection of the goodness of the fit. Two fitting approaches were used. In the first, T_i was treated as an independent fitting parameter. In the second, instead, $T_e = T_i$ was assumed which can be justified by calculating the equilibration time, τ_{ei}^e for the given set of plasma parameters (T_e, Z, n_e).

Fig. 5.17(a) shows the temperatures determined by the first approach, without making any

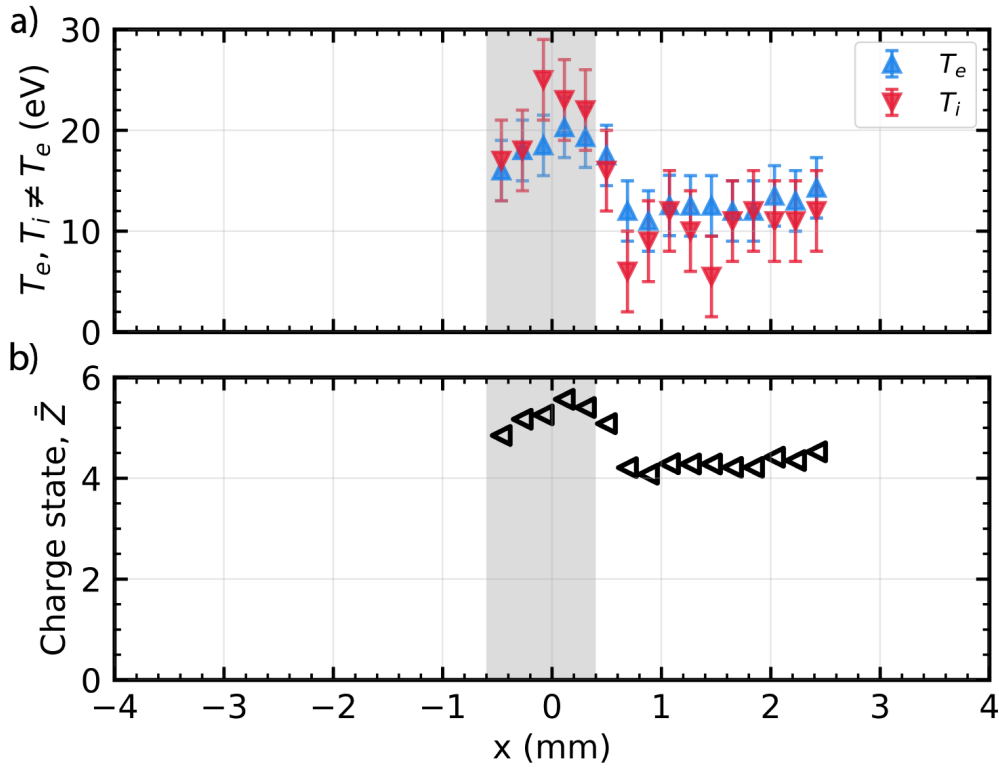


Figure 5.17: (a) Ion and electron temperatures estimated by fitting the theoretical Thomson scattering spectrum to the data, where T_i and T_e are both free parameters. The error bars are estimated by manually varying fit parameters and visually assessing the goodness of the fit. (b) Estimated average charge state, $\bar{Z}$, calculated from FLYCHK atomic code assuming the electron density n_e reported in Fig. 5.16(a) and the measured ZT_e .

assumptions on the relation between T_e and T_i . Fig. 5.17(b) shows the average ionisation state, $\bar{Z}$, derived using the determined values of T_e and the ionisation tables for silicon computed using FLYCHK atomic code. The estimated charge state by this approach is $\bar{Z} \simeq 4$ in the plasma upstream, while increases up to 5.5 within the shocked layer. Outside of the layer, the electron temperature remains relatively constant at around 12 ± 3 eV. These values are consistent with the value of temperatures obtained in previous experiments involving a similar experimental setup with single silicon targets [67], despite a different flow-magnetic field orientation. The ion temperature, T_i , follows a similar trend with the averaged measured spectral broadening corresponding to 12 eV. For the ion temperature T_i , the estimates are less accurate leading to larger error bars (± 4 eV) compared to T_e . Fig. 5.17 shows that there is a sudden increase in both temperatures, T_i and T_e , to $\simeq 18$ eV at the shock transition (fibre 10) while a higher temperature is measured at the centre of the shocked layer, reaching approximately 22 eV. This measured temperature for the ions is much lower than could be expected from shock heating if all kinetic energy of incoming ions is converted to heat without losses, which will be discussed

in more detail later in the chapter. Despite the measured T_i being slightly higher than T_e for a given spatial location within the shocked layer, the discrepancy between T_i and T_e falls well within the measurement uncertainties suggesting that the system is likely to be in thermal equilibrium with $T_e \simeq T_i$ across the entire range of measurements. This can be further supported by comparing the thermal equilibration time between electrons and ions with the characteristic time scale of the system. The plasma in this system is very collisional ($n_e \simeq 1 \times 10^{19} \text{ cm}^{-3}$, $T_e \simeq 20 \text{ eV}$, $\bar{Z} \simeq 5$) so the equilibration time, τ_{ei}^e is expected to be very short, $\simeq 2 \text{ ns}$, in comparison with the layer time scale of $\simeq 100 \text{ ns}$, hence temperature differences between ions and electrons are expected to be negligible (further details are presented later in the discussion). This lets us believe that the use of the second Thomson scattering fitting approach is well justified.

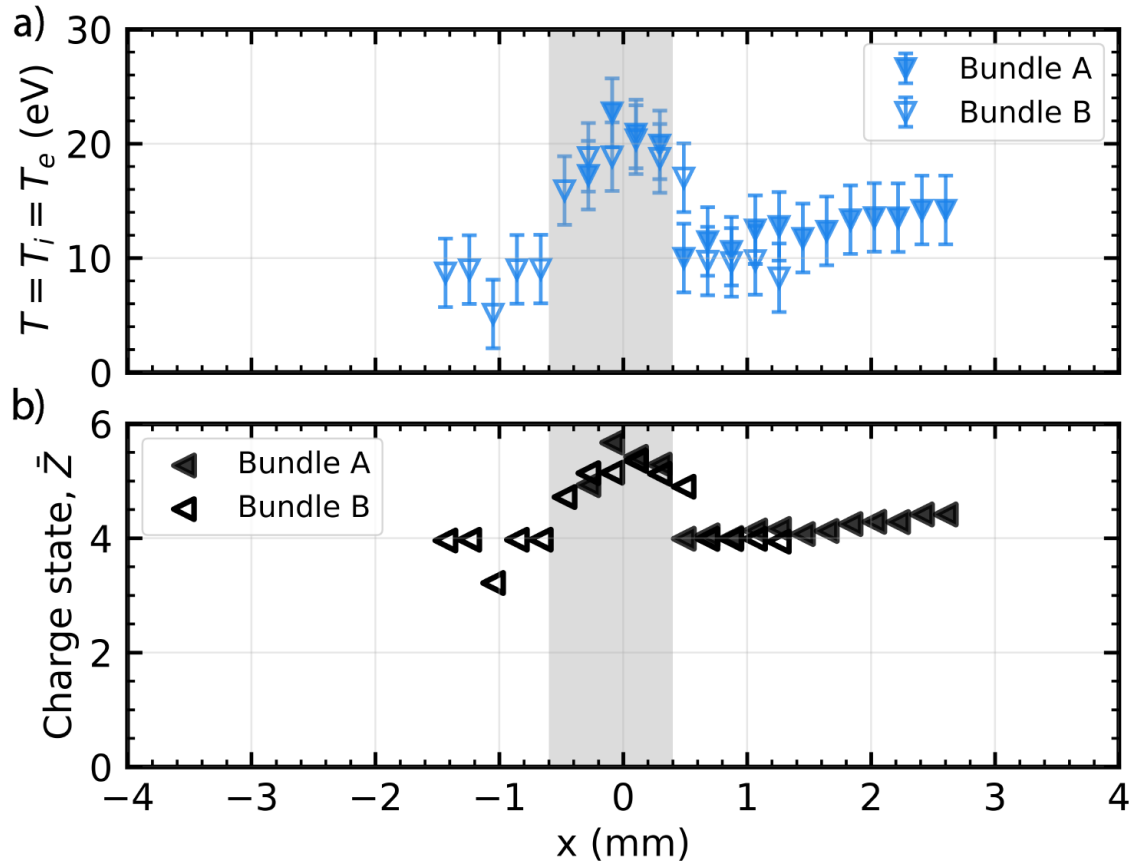


Figure 5.18: (a) Ion and electron temperatures measured for Thomson scattering spectra for both bundle A and bundle B, assuming $T_e = T_i$. The error bars are estimated by manually varying fit parameters and visually assessing the goodness of the fit. (b) Average charge state, $\bar{Z}$, calculated from FLYCHK atomic code assuming the electron density n_e reported in Fig. 5.16(a) and the measured $T_e = T_i = T$.

Results of the fitting process of the recorded Thomson scattering spectra performed assuming

$T = T_e = T_i$, are shown in Fig. 5.18(a). The average ionisation of the plasma, $\bar{Z}$, was calculated using the same atomic model and is shown in Fig. 5.18(b). Overall the fitting assuming $T = T_e = T_i$ gives values of temperatures similar to those obtained from fitting the spectra with T_e and T_i as independent parameters. Moreover, temperature measurements are also consistent across the two fibre bundles, A and B. The measured temperature, T , in the outflow slightly decreases along the flow from $T \simeq 14$ eV at $x = 2.5$ mm to $T \simeq 10$ eV just before the stagnation layer. The plasma temperature is approximately the same on the left sides of the layer of $\simeq 10$ eV. Inside the stagnated region, the plasma temperature is higher, $T \simeq 20$ eV and it appears to be slightly higher in the very centre of the layer. Although estimates of $T_i = T$ obtained in this approach are not dissimilar to those obtained with independent variations of T_e and T_i , we believe that it gives more reliable measurements of T_i .

In summary, the Thomson scattering measurements have successfully characterized both the temperature and in-plane velocity across the shock system. The velocity measurements reveal a significant reduction in velocity at the edge of the shocked layer, from $\simeq 30$ km/s to $\simeq 5$ km/s. The velocity within the shocked layer is non-zero and mostly redirected along the layer. The electron and ion temperatures show a small increase from $\simeq 12$ eV in the inflow to $\simeq 20$ eV in the stagnation layer. The small observed temperature increase in the shocked layer is lower than could be expected for shocks formed by supersonic flows. This could be due to the presence of radiative cooling effects and ionization processes which can strongly influence the plasma condition in the post-shock regions, as will be discussed in the following section.

5.1.4 Discussion of the accretion shock system

Multiple diagnostic techniques, including interferometry and Thomson scattering, have been utilised to study the shock structure resulting from the interaction of two supersonic plasma flows propagating along an ambient magnetic field. A stagnation layer is formed at collision, bounded by two reverse shocks propagating into the upstream flow. The shocked layer expands symmetrically with respect to the mid-plane as observed from the time-resolved optical self-emission images. If we consider only one side of the shocked layer, the system is effectively a free-boundary reverse shock system comparable to accretion shocks in astrophysics. In this case, the internal region of the shocked

layer represents the shocked plasma at the base of the accreting column, near the star surface, while the shocked layer boundary can instead be considered as the reverse shock front separating the upstream flow from the dense post-shock region. The ambient magnetic field is aligned to the plasma flow direction and appears to not influence the shocked structure while playing a role in collimating the two plasma streams. Thomson scattering measurements have revealed an increase in plasma temperatures and an abrupt drop in flow velocity at the position of the reverse shocks. An interesting question arises on whether the observed reverse shock evolution is in agreement with the 1-D shock model, normally assumed to interpret astrophysical observations. In this context, the position of the shock front can have important consequences for the explanation of the luminosity curve observed in stellar systems, thus benchmark accretion models in the lab can lead to a better understanding of the star formation mechanism.

In this section, we present a detailed discussion of the key findings derived from the data analysis by comparing the system with the 1-D shock model. A scaling comparison with the astrophysical case is then performed, focusing on the expected reverse shock position. The significance of radiative cooling effects in the post-shock region is also addressed. Finally, the discussion concludes with several open questions that highlight areas for future research and investigation.

Properties of the reverse shock system

In the previous sections, hydrodynamic parameters such as density, temperatures and velocity were successfully characterised in both the shocked layer and the plasma inflows. It was found that the plasma inflows were supersonic, directed along the target normal, with a velocity of $\simeq 30$ km/s just upstream of the stagnation layer, density of $\simeq 5 \times 10^{18} \text{ cm}^{-3}$ and electron and ion temperatures of $\simeq 12$ eV. In the stagnated plasma, behind the two reverse shocks, both temperature and density are higher, while velocity is reduced to become below the characteristic sonic speed.

A summary of the measured plasma parameters is presented in Table 5.1, including both the inflow plasma conditions and the properties of the shocked layer. It is important to acknowledge that these values were determined at a specific time point, approximately 220 ns after the x-ray pulse. Consequently, the upstream and downstream conditions may vary throughout the course of the

experiment.

Quantity	Symbol	Inflow (at $x = \pm 0.5$ mm)	Shocked layer (at $x = 0$ mm)
Flow velocity (km/s)	V_x (V_y)	30 ± 1	(5 ± 1)
Electron density ($\times 10^{19} \text{ cm}^{-3}$)	n_e	0.4	1.5
Temperature (eV)	$T = T_e \simeq T_i$	12 ± 3	20 ± 3
Average ionization	$\bar{Z}$	4	5
Ambient B-Field (T) (estimated)	B	$\simeq 5$	$\simeq 5$

Table 5.1: Characteristic plasma parameters of silicon flow at $\simeq 220$ ns after x-ray pulse start. $\bar{Z}$ is computed using FLYCHK code [52] assuming measured T_e and interpolating the available ionization curves to better approximate the measured electron density, n_e .

In this configuration, the ambient magnetic field is parallel to the inflow direction, which means that no B-field compression is expected with no direct effect on the shock jump condition. The ambient field, however, can contribute to the collimation of the plasma streams and potentially affect the plasma expansion from the shocked layer in the transverse direction. An estimate of the ambient magnetic field at the layer position at distance $r = 28$ mm from the Z-pinch and at the time of plasma formation, $\simeq 300$ ns after the current pulse, can be calculated from the current profiles presented in Fig. 5.3 (a) as:

$$B(300 \text{ ns}) = \frac{\mu_0 I(300 \text{ ns})}{2\pi r} \simeq 5 \text{ T} \quad (5.1)$$

This estimate is consistent with previous b-dot measurements in cylindrical wire array experiments and corresponds to a magnetic pressure $P_B = B^2/2\mu_0 \simeq 10 \text{ MPa}$ directed perpendicularly to the streaming plasma flows. Thermal properties of the plasma flows are the dominant factors that set the transverse divergence of the plasma flow. In this case, effective collimation by the magnetic field would happen when magnetic pressure is much higher than the plasma thermal pressure so that thermal beta $\beta_{th} \ll 1$. In the upstream plasmas, denoted with subscript 1, the thermal pressure is defined as $P_{th,1} = n_{i,1} k_B (Z_1 + 1) T_1 \simeq 9.5 \text{ MPa}$ for $T_1 = T_e = T_i \simeq 12 \text{ eV}$ and $n_{i,1} \simeq 5 \times 10^{18} \text{ cm}^{-3}$, hence thermal beta, $\beta_{th} = P_{th}/P_B = 0.95$. Under this condition, the ambient magnetic field can only partially contribute to the overall collimation of the upstream plasma flows. The two dominant pressure terms in the reverse shock dynamics are therefore the thermal pressure of the layer and the ram pressure of the upstream flows. These terms are expected to be balanced at the reverse shock. Using the upstream plasma parameters reported in Table 5.1, the ram pressure of the inflows is given by $P_{ram,1} = n_{i,1} m_i V_1^2 \simeq 42 \text{ MPa}$. Within the stagnation layer, thermal pressure is dominant with

$P_{th,2} \simeq 50$ MPa which is reasonably close to the pressure inflow, $P_{th,2}/P_{ram,1} \simeq 1.2$. These pressure contributions are also much larger than the magnetic pressure, P_B , almost 5-fold larger, leading to high-beta conditions both in the upstream, $\beta_{dyn} \simeq 4$ and in the shocked layer, $\beta_{th} \simeq 5$. Thus, the magnetic field is not sufficient to prevent the transverse expansion of stagnated plasmas.

In astrophysical accreting columns, the dominant β can vary over a very broad range from $\beta \ll 1$ in high-field systems such as MCVs (magnetic cataclysmic variable stars) [111, 112], to $\beta \gg 1$ in high-beta T Tauri stars or Herbig Ae/Be. The astrophysical simulations presented in [26], investigated the properties of accretion flows, specifically for cases with $\beta \gg 1$ in the stagnation region, looking at radiative cooling effects. It is important to highlight that these conditions can dynamically change as magnetic fields and plasma properties of the accreting columns evolve over time. In these systems, conditions of high β_{th} are likely to occur at the early formation of the post-shock region while reducing as more outflows from the stagnation are radially escaping and compressing the surrounding magnetic field. In the laboratory experiments presented here, the geometry of the inflow and of the stagnation layer means that for the central part of the shocked layer, corresponding to the target width, $\simeq 5$ mm, the interaction is one-dimensional and the properties of the process can be well compared with the 1-D models and numerical simulations in [26]. We, however, did not observe any effects of the magnetic field on the transverse expansion of the layer (formation of additional compression in the transverse outflow as in [26]). One possibility relies on the large value of β estimated in the experiment, but it is also possible that the magnetic field is not well frozen in the plasma at the spatial scales in the experiment. In the shocked layer, magnetic diffusion becomes important at the spatial scale given by $Re_m = c_{s,2}L_\eta/D_M = 1$, which corresponds to $L_\eta \simeq 0.87$ mm. Despite these differences in the post-shock region between the astrophysical systems and the experiments, in both cases, the magnetic field plays a marginal role in the dynamic of the reverse shocks so instructive comparisons can still be made assuming a 1-D accretion shock model.

Comparison with a 1-D reverse shock model

In its simple form, a reverse shock system is formed when supersonic plasmas collide against a hard surface, which provides a supporting pressure for the stagnant plasma, as presented in Chapter 4. In our experiment with the collision of two identical flows, the role of the stationary obstacle is

replaced by the mid-plane of the stagnation layer, which remains stationary due to the presence of two equal and opposite ram pressures of the flows. This leads to the formation of two reverse shocks and a common post-shock region. By considering only half of the shocked structure, the reverse shock system is equivalent to the single flow case, so that the observations can be benchmarked with a classical shock model. The system is symmetric, therefore considerations made on one side of the shocked layer are still valid for the opposite reverse shock.

While shock waves are entropy-increasing processes that convert kinetic energy into heat, the continuity of plasma properties (jump conditions) is maintained across shocks as long as the upstream and downstream conditions are sufficiently far from the shock front. In this analysis, the shock transition is effectively treated as a one-dimensional discontinuity, collapsing into a narrow region of negligible thickness. The properties of the system are described by the Rankine-Hugoniot relations [23, 25], which will be utilised to characterise the strength and properties of the observed reverse shock. In unmagnetized plasmas, where the magnetic field is not dynamically significant (or for flow parallel to the magnetic field), the behaviour is analogous to classic fluid dynamics, and the characteristic speed of the shock waves is determined by the speed of sound, denoted as c_s . When two plasma flows interact with each other at speeds greater than the speed of sound, the system cannot instantaneously dissipate the resulting pressure changes, leading to the formation of a shock wave. For this class of shocks, the dominant dissipation mechanism is the ion viscous effect, through ion-ion Coulomb collisions. The mean free paths are typically very small in highly dense plasmas, indicating a highly collisional system. The shock transition length is comparable to the ion mean free path calculated with respect to the inflow velocity $\lambda_{ii} = V/v_i$:

$$\lambda_{ii} = 4 \times 10^8 \frac{VT_i[eV]^{3/2}}{Z^3 A^{-0.5} n_e \Lambda} \simeq 0.04 \mu m \quad (5.2)$$

Here, $A = 28$ is the atomic weight of silicon, n_e is the electron density in cm^{-3} , Z is the charge state, and $\Lambda \simeq 4$ is the Coulomb logarithm. For completeness, all relevant characteristic lengths which in various situations could affect the properties of the shock transition, are summarised in Table 5.2 [18, 113, 114].

In our experiments, it is not possible to resolve spatial scales of the order of the ion mean free path,

Quantity	Symbol	Inflow (at $x = \pm 1$ mm)
Ion-ion mean free path (μm)	λ_{ii}	0.04
Electron-ion mean free path (μm)	λ_{ei}	0.003
Resistive Diffusion Length (μm)	L_η	870

Table 5.2: Characteristic spatial scales calculated at $\simeq 220$ ns after the X-ray pulse start. The magnetic diffusivity is estimated assuming Spitzer resistivity, η , [104], given by $D_M = c^2\eta/4\pi$.

λ_{ii} , however, an upper limit of the shock transition length can be estimated by examining line-integrated interferometry measurements, as demonstrated in Fig. 5.9, and the horizontal density profiles. This length corresponds to the spatial extent over which the electron density increases from its pre-shock value to the post-shock value. Based on these considerations, the estimated shock transition length is less than $200 \mu\text{m}$, however, this represents the upper limit of the transition due to line integration and alignment/front uniformity effects along the z -axis. This upper limit estimation is consistent with the upper limit seen from Thomson scattering measurements, where changes in velocity and temperatures at the boundary of the shocked layer occur at scales smaller than the distance between scattering volumes (approximately $200 \mu\text{m}$). In this configuration with plasma flow parallel to the magnetic field, we do not expect any currents at the shock front, in contrast to the situation where the magnetic field is perpendicular to the flow, leading to some B-field compression and hence corresponding current. This rules out the possibility of shocks mediated by resistive dissipation ($L_\eta \simeq 0.87$ mm, Table 5.2) and indicates that the shock transition is primarily mediated by Coulomb collisions, similar to conventional hydrodynamic shocks with characteristic transitions of a few mean free paths. For the measured inflow plasma parameters, the shortest collision mean free path is represented by the electron-ion mean free path, $\lambda_{ei} = V/v_e \simeq 0.003 \mu\text{m}$, one order of magnitude lower than λ_{ii} . This difference between the two characteristic mean free paths arises from the relatively high averaged charge state, $\bar{Z} \simeq 5$ of the silicon plasma which can have significant implications on the thermal diffusivity across the shock as it will be discussed later in the chapter.

For the following discussion of the measurement of the reverse shock properties, a schematic representation of the shock system is illustrated in Fig. 5.19 for the lab and the shock frame of reference. Two regions can be distinguished across the reverse shock front, an upstream or pre-shock region (subscript 1) and a downstream or post-shock region (subscript 2), representing conditions in the plasma inflow and in the shocked layer, respectively. In the shock frame of reference, the shock

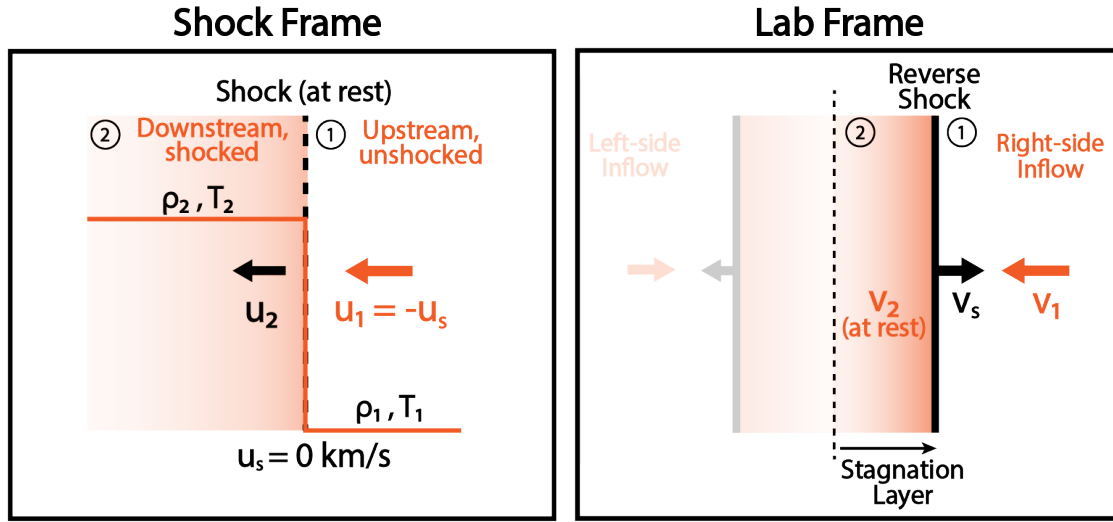


Figure 5.19: (Left) Schematic representation of a one-dimensional shock in the shock frame of reference. The shock is at rest with the post-shocked flow leaving the shock front at velocity $u_2 < u_1$. (Right) Reverse shock in the laboratory frame where the shock front is moving at speed v_s into the upstream flow, while the downstream flow is at rest, $v_2 = 0$ km/s.

front is at rest with the incoming flow moving at u_2 and post-shock flow with $u_2 < u_1$. The shock relations are typically derived and present for this frame of reference, as illustrated in Fig. 5.19(Left). In the situation when the post-shock plasma is not moving supported by a rigid unmoving obstacle, or by a second identical flow, the shock front cannot be stationary due to material accumulation. This is equivalent to considering a reference frame moving with velocity u_2 by subtracting from all the velocities the upstream velocity, u_1 . Thus, the shock front in this reference frame, equivalent to the lab frame in the experiment, will be moving with velocity $v_s = -u_2$ in the direction opposite to the incoming flow, as shown in Fig. 5.19(Right). The relations between the experimental measurements of v_s , v_1 and the equivalent parameters in the shock frame of reference are given by:

$$\begin{aligned} u_1 &= v_1 + u_2 \\ u_2 &= -v_s \\ v_2 &= 0 \text{ km/s} \end{aligned} \tag{5.3}$$

When analysing shock systems, it is useful to express the relations presented in (5.3) in terms of ratios of density and velocity across the shock front defined as $R = \rho_2/\rho_1 = u_1/u_2$. In the strong shock limit, this shock jump ratio depends solely on the polytropic index γ . By substituting R into the equations

(5.3), the velocities are then written as follows:

$$\begin{aligned} \mathbf{u}_1 &= \mathbf{v}_1 \frac{R-1}{R} \\ \mathbf{v}_s &= -\mathbf{u}_2 = -\frac{\mathbf{v}_1}{R-1} \end{aligned} \quad (5.4)$$

From these considerations, dimensionless parameters can also be defined. In particular the upstream sonic Mach number in the shock frame, $M_{u,1} = u_1/c_{s,1}$, where $c_{s,1}$ is the characteristic sonic speed. In this case, the sound speed can be defined as:

$$c_{s,\alpha} = \sqrt{\frac{k_B (\bar{Z}_\alpha T_{e,\alpha} + T_{i,\alpha})}{m_i}} \quad (5.5)$$

α indicates the two regions under consideration: the pre-shock and post-shock regions. As shown in Sec. 5, it was revealed that the shocked layer is not stationary but rather expands towards the upstream direction with a speed of 13 ± 0.5 km/s, corresponding to a measured reverse shock speed of $v_s \simeq 7$ km/s. If we now move to the shock frame using eq. 5.3 and considering the shock front at rest, the upstream speed is given by $u_1 = v_1 + v_s \simeq 37$ km/s, which differs insignificantly from the measured upstream velocity in the lab frame. Based on the measured plasma parameters reported in Table 5.1, the sound speeds in the plasma flows and the shocked layer are $c_{s,1} \simeq 14.5$ km/s and $c_{s,2} \simeq 21$ km/s, respectively. The calculated upstream Mach number is $M_{u,1} \simeq 2.5$. For the model in Fig. 5.19(Left), the post-shock velocity, v_2 , should be zero, however, Thomson scattering measurements revealed a non-zero velocity within the shocked layer, $v_2 \simeq 5$ km/s. This fact could be just associated with deviations of the Thomson scattering geometry from expected or with internal dynamics in the layer, despite these having symmetry and thus, $v_2 = 0$ km/s at the centre. An expression commonly used to estimate the expected density jump is given by:

$$\frac{\rho_2}{\rho_1} = \frac{M_u^2(\gamma+1)}{M_u^2(\gamma-1)+2} \quad (5.6)$$

Fig. 5.20 illustrates the expected compression ratio as a function of the polytropic index, γ , for various upstream Mach numbers. The rectangular region in the figure represents the range of the measured compression ratios from the interferometry data. The measured values of R_{n_e} discussed in § 5.1.2, are

adjusted by taking into account the ionisation factor across the shock using $\bar{Z}_1$ and $\bar{Z}_2$ from Thomson scattering fitting, as follows:

$$R = \frac{\bar{Z}_1}{\bar{Z}_2} R_{ne} = \frac{4}{5} R_{ne} \quad (5.7)$$

The shaded region indicates the confidence level of the measurements. In the case of large upstream Mach numbers, Equation 5.6 is simplified to the strong shock limit relation, which is solely dependent on γ . From Fig. 5.20 it is seen that the measured compression ratio, $R = 4 \pm 0.5$, is not compatible

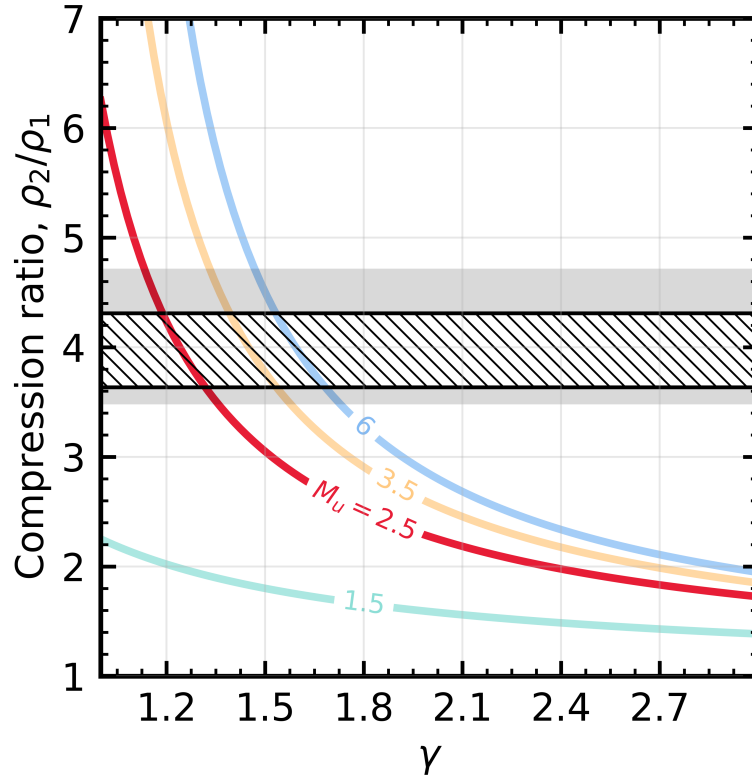


Figure 5.20: Compression ratio, ρ_2/ρ_1 as a function of the polytropic index, γ for a given upstream Mach number, M_u .

with the assumption of an ideal gas with $\gamma = 5/3$ for an upstream Mach number, $M_{u,1} = 2.5$. In order to match the compression ratio measured from the electron density map, a value of $\gamma \leq 1.2$ is required. This suggests that both ionisation and radiative cooling effects can potentially play a significant role in the reverse shock dynamic, leading to a low effective γ . The temperature difference across the shock front recorded by Thomson scattering is enough to increase the average charge state of the plasma, $\bar{Z}$, from 4 to 5, suggesting that ionisation processes are indeed not negligible. The value of γ significantly influences the compressibility of the plasma, thereby affecting shock transition lengths and the ratios of plasma parameters across the shock front. Additionally, it impacts significantly the reverse shock

velocity, which represents one of the main parameters for astrophysical observations.

To check that the value of $\gamma \leq 1.2$ is reasonable in this plasma system, a theoretical estimate of the polytropic index is given by the relationship $\gamma = 1 + p/(\rho\epsilon)$, where p denotes pressure and $\rho\epsilon$ represents the energy density [25]. For simplicity, only a single value of the polytropic index is considered between the upstream and post-shock regions. Assuming $T_e \simeq T_i \simeq 18$ eV in the ionizing regime (not fully stripped), the expressions for pressure and energy density in CGS units are as follows:

$$p = 1.6 \times 10^{-12} n_i \left(T_e(1 + \bar{Z}) - 6.96 \times 10^{-8} n_i^{1/3} \bar{Z}^2 \right) \quad (5.8)$$

$$\rho\epsilon = 2.4 \times 10^{-12} n_i \left(825 + T_e(1 + \bar{Z}) - 1.39 \times 10^{-7} n_i^{1/3} \bar{Z}^2 \right) \quad (5.9)$$

By substituting the measured plasma parameters within the shocked layer ($\bar{Z} \simeq 5$ and $n_i \simeq 0.3 \times 10^{19} \text{ cm}^{-3}$), we find that the polytropic index is approximately $\gamma \simeq 1.1$. This is consistent with the expected value of γ for the experimentally measured upstream Mach number, $M_{u,1}$ and compression ratio, R , from the relation, (5.6) shown in Fig. 5.20. It is important to note that equations (5.8) and (5.9) take only into account ionisation and the presented model is not very sensitive to both temperature and charge state, therefore we proceed with the discussion by considering the experimental estimate of γ .

Carrying on the analysis with $\gamma = 1.2$, the predicted shock velocity and temperature jumps can be also compared. The expected velocity jump $u_1/u_2 = 1/R \simeq 0.3$ is consistent with the observations from Thomson scattering where the measured velocity ratio across the shock is $\simeq 0.2$. Using the equations (5.4) and a compression ratio $R = 4.3$ the estimated shock velocity is $v_s = -11$ km/s slightly higher than the measured value, $\simeq -7$ km/s. A possible explanation for this can be attributed to the fact that the experiment is not completely 1-D as it is observed from the imaging data. The shocked layer extends beyond the silicon target size, perpendicular to the surface normal, with a decreasing density distribution along the layer. This, in combination with a non-zero velocity within the post-shock region from Thomson scattering measurements, indicates that transverse outflows leave the central stagnation region, moving towards the edges of the layer, without contributing to the shock accretion. Also, the uncertainties associated with shocked layer width from optical self-emission images due to both line-integration effects and low spatial resolution can have a significant impact on the experimental

estimate of the shock velocity, therefore it is reasonable to state that the shock velocity is also consistent with the 1-D accretion model.

A schematic representation of the expected temperature profiles at a shock front is illustrated in Fig. 5.21. Most of the kinetic energy in the incoming flow is carried by ions due to their large mass, so they are the ones heated at the shock transition by viscous effects, while the kinetic energy of the incident electrons is only negligibly converted into electron heat. The temperature difference of the electrons is smaller by a factor of m_e/m_i than ΔT_i and it typically happens through adiabatic compression. In the upstream plasma flow ions and electrons are "rigidly coupled" with a single

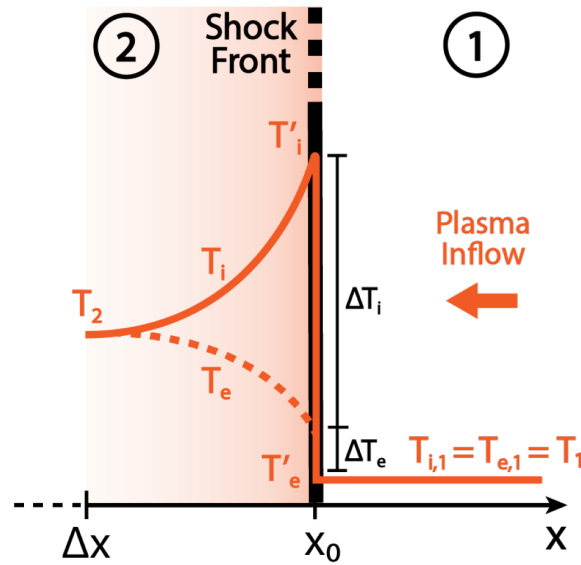


Figure 5.21: Ion and electron (dashed line) temperature profiles in a plasma shock front without any thermal heat conductions.

temperature, T_1 due to high collisionality. At the shock front, ions and electrons are instantaneously heated at temperatures T'_i and T'_e respectively, representing the maximum available thermal energy after shock compression. These are also the final temperatures of the post-shock region in the absence of any cooling or exchange effects between ions and electrons. If energy exchange occurs through Coulomb collision, a relaxation region is formed with length Δx where T_i and T_e equilibrate until reaching the final post-shock temperature, T_2 . The time of equilibration, τ_{ei} , depends inversely on the collision rate of ions and electrons and sets the distance from the shock front within which the two temperatures are distinguishable. In § 5.1.3, it is shown that differences between ion and electron temperatures are negligible for all positions in the layer and well within the range of uncertainties.

This raises two natural questions: Should we expect any temperature differences between ions and electrons? And if yes, is this discrepancy large enough to be detected at the spatial resolution capability of the Thomson scattering diagnostic (distance between two fibres, $\simeq 200 \mu\text{m}$)?

To address these questions, we start by estimating the thermal equilibration time between ions and electrons, assuming fully thermalised ions and cold electrons. The recorded temperatures of the incoming flow are $T_1 = 12 \text{ eV}$ for both ions and electrons. To estimate the maximum increase of temperatures at the shock front, a rough estimate is derived from the kinetic energy of the flow with $V_x \simeq 30 \text{ km/s}$. The ion temperature jumps are given by:

$$\Delta T_i [\text{eV}] = \frac{2}{3} E_i = \frac{2}{3} \frac{m_i V_x^2}{2e} \simeq 90 \pm 6 \text{ eV} \quad (5.10)$$

The increase in electron temperature is negligible due to their small mass, $\Delta T_e \simeq 0 \text{ eV}$. In reality, the electron temperature increases due to adiabatic compression, $T'_e = T_1 R^{\gamma-1} \simeq 16 \text{ eV}$ for $R = 4$ and $\gamma = 1.2$, however, this increase falls within the range of temperature uncertainties so it is reasonable to assume $T'_e \simeq T_1$. If any energy exchanges between ions and electrons are neglected, the maximum temperatures are therefore estimated to be $T'_i \simeq 102 \pm 6 \text{ eV}$ and $T'_e = 12 \text{ eV}$, knowing that $\Delta T_i = T'_i - T_1$.

The thermal equilibration time for both species can be calculated using their respective collision frequencies as given by [113]:

$$\tau_{\alpha\beta}^e = \frac{1}{\bar{v}_e^{\alpha|\beta}} \quad (5.11)$$

$$\bar{v}_e^{\alpha|\beta} = 1.8 \times 10^{-19} \frac{(m_\alpha m_\beta)^{1/2} \bar{Z}^2 n_\beta \Lambda}{(m_\alpha T_\beta + m_\beta T_\alpha)^{3/2}} \quad (5.12)$$

Here, α and β represent generic plasma species such as ions and electrons. We assume a single average charge state in the post-shock region, $\bar{Z} \simeq 5$, and for this calculation, the Coulomb logarithm Λ is approximately 4. By substituting the post-shocked region's plasma temperatures (in eV) and densities (in cm^{-3}), the calculated thermal equilibration time is $\tau^e \simeq 1 \text{ ns}$ [114]. This short equilibration time suggests that immediately after the shock transition, ions and electrons are expected to be in thermal equilibrium with $T_i \simeq T_e$. At this timescale ($\simeq 1 \text{ ns}$) and plasma velocities ($\simeq 5 \text{ km/s}$), the spatial scale on which T_e and T_i equilibrate is very short, $\Delta x \sim u_2 \tau^e \simeq 3.5 \mu\text{m}$ and would not be resolvable with

the Thomson scattering diagnostic, $\Delta x \ll 200 \mu\text{m}$.

Within the shocked layer, the measured temperatures are $\simeq 18 \pm 3 \text{ eV}$ at the boundaries of the shocked layer, $x = \pm 0.5 \text{ mm}$ and $\simeq 24 \pm 3 \text{ eV}$ at the layer centre, which is more than 4 times lower than the maximum estimated ion temperature, T_i' . Until now we have considered the hypothesis of full thermalisation so that the entire available kinetic energy of the plasma is efficiently converted into heat, neglecting any losses. This is a good approximation in strong shock regimes when the upstream Mach number is infinitely large, $M_u \rightarrow \infty$. This, however, is not the case in the experiment, where $M_u \simeq 2.5$.

A more realistic estimate of the final temperature in the post-shock region can be derived from the definition of thermal pressure in the shocked layer, $P_{th,2} = P_2 = n_{i,2}k_B(Z_2T_{e,2} + T_{i,2})$. The expected pressure ratio across the shock is calculated by the following relation:

$$\frac{P_2}{P_1} = \frac{2M_u^2\gamma - (\gamma - 1)}{\gamma + 1} \quad (5.13)$$

The corresponding pressure curves for various values of the upstream Mach numbers are shown in Fig. 5.20(Right). Again for the combination of γ value and $M_{u,1}$ the expected pressure ratio is $P_2/P_1 = 6$. Knowing the pressure ratio and the thermal pressure of the upstream plasma, $P_1 = P_{th,1} = 9.5 \text{ MPa}$, the post-shocked pressure is $P_2 = 6P_1 \simeq 57 \text{ MPa}$ and thus the estimated final temperature, T_2 :

$$T_2[\text{eV}] = \frac{Z_2P_2}{e(Z_2 + 1)n_{e,2}} \simeq 20 \text{ eV} \quad (5.14)$$

This value of post-shock temperature is remarkably close to the one observed experimentally for the final temperature at the equilibrium. Although the finite upstream Mach number can explain the detected low downstream temperature, it is important to highlight that this calculation is valid only far enough from the shock front into the shocked layer, where both plasma species are at equilibrium without necessarily representing the temperature jump at the shock front, T_i' . At one fibre distance, $\simeq 200 \mu\text{m}$, across the shock front, electrons and ions can be already considered at equilibrium ($\Delta x \simeq 3.5 \mu\text{m}$) therefore the analysis with $\gamma = 1.2$ is consistent with the measured temperatures from the first two available fibres within the layer (fibres 6 and 11).

If we repeat the analysis for $\gamma = 1.1$, the results follow better the experimental measurements,

particularly for the density and velocity ratios across the shock. The comparison is presented in Table 5.3. The predicted reverse shock velocity remains always higher than the measured value of v_s , even for values of the polytropic index as low as $\gamma = 1.1$. This can be attributed to the three-dimensional nature of the shocked layer. Although the central part of the stagnant region extending over one target width is reasonably one-dimensional, any transverse outflows along the shocked layer can significantly contribute to reducing the reverse shock expansion. Under these conditions, a factor of two in the reverse shock estimates is not surprising.

Quantity	Symbol	Measured	$\gamma = 5/3$	$\gamma = 1.2$	$\gamma = 1.1$
Density ratio	$R = \rho_2/\rho_1$	4 ± 0.5	2.8	4.3	4.7
Electron density ratio	$R_{ne} = RZ_2/Z_1$	5 ± 0.5	3.5	5.4	5.9
Velocity ratio	u_1/u_2	0.2 ± 0.04	0.4	0.23	0.2
Pressure ratio	P_2/P_1	5 ± 1.5	7.5	6	6
Shock velocity (km/s)	v_s	-7	-20.5	-11	-10
Postshock temperature (eV)	T_2	18 ± 3	25	20	20

Table 5.3: Comparison of the characteristic parameter ratios across the reverse shock, between observation and 1-D shock model with various values of γ .

Summary of comparison with 1-D shock model

Direct and indirect measurements of the plasma parameters upstream and downstream of the shock structure revealed good consistency with a 1-D accretion shock model. The key findings from the conducted analysis are summarised below:

- The collision between the two plasma flows occurs at the mid-plane with an upstream velocity just before the interaction of $V_1 \simeq 30 \pm 1$ km/s, corresponding to an upstream Mach number of $M_{u,1} \simeq 2.5$. The interaction is symmetric with respect to the mid-plane perpendicular to the flow direction, which results in the formation of a stagnation layer bounded by two reverse shocks.
- The ambient magnetic field does not affect the dynamic of the shocked layer due to the large and dominant $\beta_{th} \simeq 5$ within the layer. The orientation of the field lines (parallel to the flow direction) can marginally help to collimate the two plasma inflows.
- The density and temperatures increase within the shocked layer. The measured compression ratio is $R \simeq 4 \pm 0.5$. Both ion and electron temperatures are approximately 30% higher than the

plasma inflow, ranging from $T_1 \simeq 12 \pm 3$ eV to $T_2 \simeq 18 \pm 3$ eV just after the shock front. The short thermal equilibration time scale of approximately 2 ns justifies the absence of significant temperature differences between species.

- A comparison with the 1-D shock model has been performed, revealing consistency with the measured plasma parameters across the shocked layer for values of the polytropic index, $\gamma < 1.2$. This suggests that both ionization and cooling effects may contribute to the dynamic of the observed shock system.
- Two reverse shocks are formed at approximately 90 ns after the x-ray pulse, propagating into the upstream with a speed of $v_s \simeq 7$ km/s. The expected shock velocity is $\simeq -10$ km/s for $\gamma = 1.1$, more than twice higher than measured from optical self-emission data, likely due to transverse outflows escaping the shocked layer. The flow in the post-shock region is subsonic with a sonic Mach number of $M_{S,2} \simeq 0.34$.

Evidence of shock precursors due to radiative effects

The temperature profiles illustrated in Fig. 5.21 do not take into consideration the important effect of electron heat conduction on the shock front structure or any radiative effects. As already mentioned in the previous § 5.1.2, experimental interferometry data (Fig. 5.9 and Fig. 5.12) revealed an increase in density ahead of the shock transition, resembling a preheating precursor, Fig. 5.22. These features can be explained by possible heat transfer from the shock front to the upstream flow due to electron thermal conduction or radiative cooling effects. Until now we have assumed that dissipative processes, such as viscosity and heat conduction, occur only in the post-shock regions. Although often neglected, the effect of heat conduction can be significant in both the post-shock relaxation region and ahead of the shock front. The characteristic scale of heat conduction effects is dependent on transport coefficients and shock velocity. Thermal diffusivity, in mid-Z plasmas, is determined by the electron thermal diffusivity since the electron mean free path, λ_e is an order of magnitude higher than the ion mean free path, λ_i , and the electron velocity is larger by a factor of $(m_i/m_e)^{1/2}$ than the respective values for the ions. Therefore the characteristic scale over which the electron heat conduction takes place is

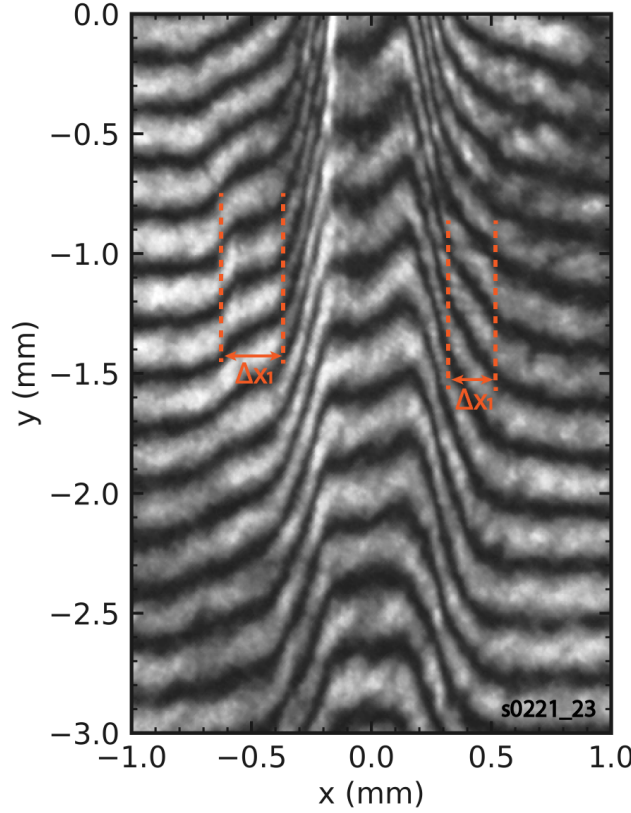


Figure 5.22: Zoom-in region of the shocked layer taken from the early time shot (s0221_23) at 174 ns after the X-ray pulse presented in Fig. 5.12. A precursor region is observed on both sides

$l_e \sim \chi_e/u_1 \simeq 30 \text{ } \mu\text{m}$ where the electron thermal diffusivity is given as, $\chi_e(T_{e,2}) = k_{||}/n_{e,2} \simeq 9.3 \times 10^3 \text{ cm}^2/\text{s}$ with the conductivity of the electrons is $k_{||} = k^c k_B n_{e,2} T_{e,2} \tau_e / m_e$ by [104] (see Table 5.5).

Fig. 5.23 shows temperature profiles across the shock when electron thermal conduction is considered [23]. If we assumed x_0 at the point of the shock front, the thickness of the preheating layer ahead of the shock front is of the same order as the thickness of the relaxation region behind the shock, $|\Delta x_1| \simeq \Delta x$. Therefore, the thickness of the pre-heating region is dependent on the electron thermal diffusivity estimated assuming the downstream temperature T_2 at equilibrium, $\chi_e(T_2)$, given by ([23], eq. 7.36):

$$|\Delta x_1| = \frac{2}{5} \frac{\chi_e(T_2)}{u_1} \simeq 25 \text{ } \mu\text{m} \quad (5.15)$$

Where the upstream velocity is $u_1 = 37 \text{ km/s}$. This estimated thickness of the pre-heating region due to electron thermal conduction is very small compared to the precursor length observed from the raw interferograms, $\simeq 250 \text{ } \mu\text{m}$, indicating that thermal conduction cannot justify these density features. An alternative mechanism can be represented by heat transfer due to radiative absorption. Radiation

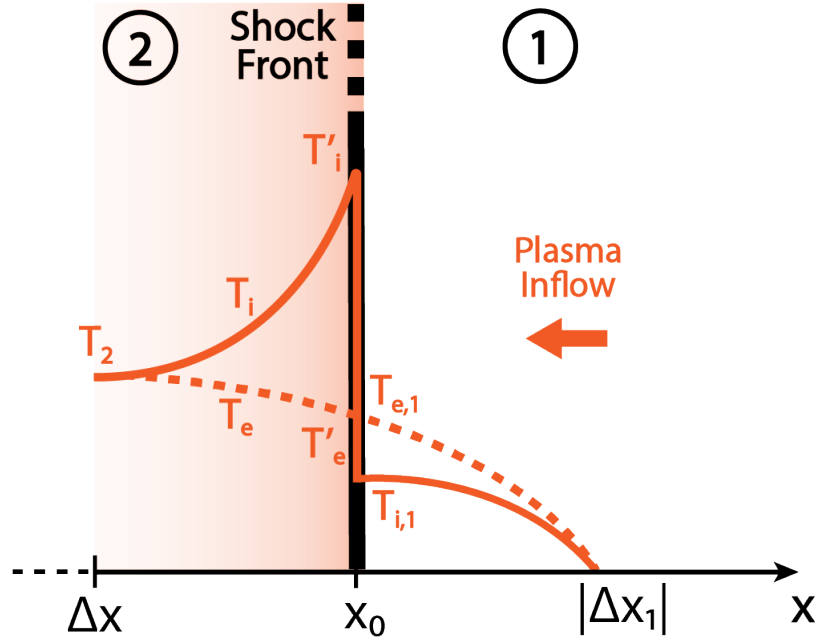


Figure 5.23: Ion and electron (dashed line) temperature profiles in a plasma shock front without any thermal heat conductions.

from the shocked heated material travels upstream to heat the incoming flow before reaching the shock front.

To test this hypothesis, numerical simulations of the shock system were carried out using the 1-D radiation-magnetohydrodynamics code HELIOS-CR [68]. The simulations were initiated with two equal-sized regions of silicon plasma at constant $\rho_0 = 4.7 \times 10^{-5} \text{ g/cm}^3$ and temperatures $T_e = T_i = 10 \text{ eV}$. The flow velocity of $u_1 = 37 \text{ km/s}$ was equal in magnitude for both plasma flows but with opposite directions, such that a stagnation layer was formed at the centre of the domain with two reverse shocks propagating upstream, similar to the experiments. The plasma flow was free to expand at the boundaries so that the propagation of the reverse shocks was supported by a constant plasma inflow. Results of these simulations at the time when the shock has reached a steady-state phase ($\simeq 40 \text{ ns}$) are shown in Fig. 5.24. The system is symmetric with respect to the boundary of the two regions and only the reverse shock propagating from left to right is presented in the figures.

Fig. 5.24(a) shows the shock structure in the case of only thermal conduction and no radiation transport. Compression, heating and ionisation occur at the shock front ($x = 1 \text{ mm}$) with ion temperature $T_i = 20 \text{ eV}$ and density jump $\rho_2/\rho_1 = 5$, consistent with the observations. The mean charge state

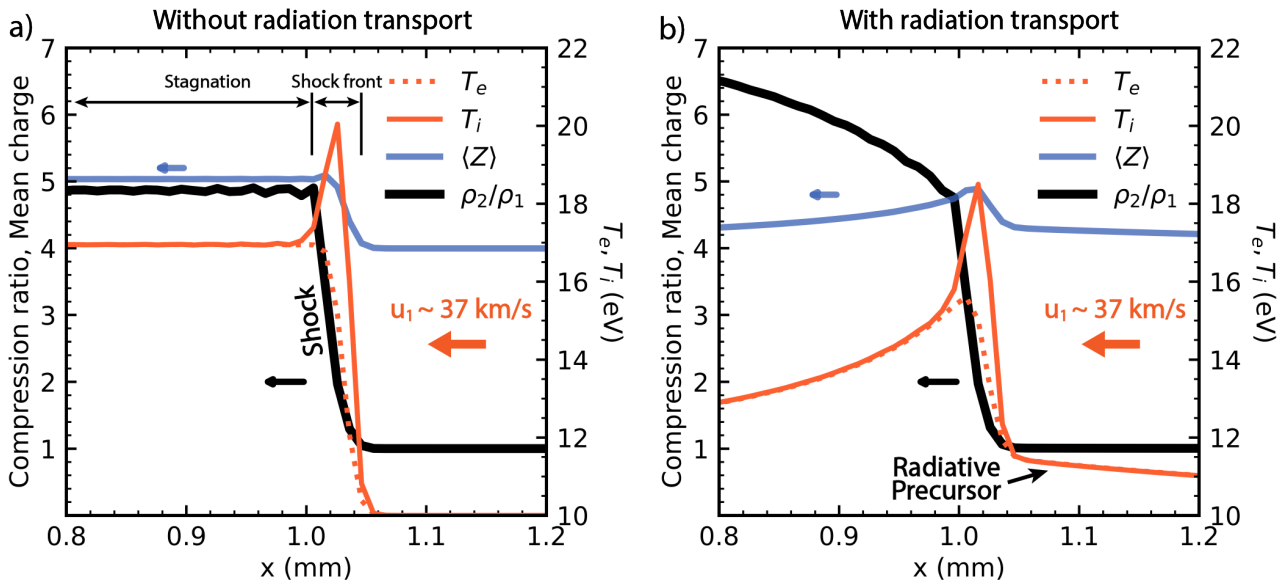


Figure 5.24: HELIOS-CR [68] simulation of reverse shocks formed from the collision of counter-streaming silicon plasma flows. The system is symmetric so only one side of the shocked layer is shown in the figure. The shock moves from left to right supported by an upstream velocity, $u_1 = 37$ km/s. (a) Compression, heating and ionisation profiles across the shock assuming only thermal conduction and no radiation losses. No preheat region is seen ahead of the shock. (b) Compression, heating and ionisation profiles across the shock with radiation transport switched on. A radiative precursor is observed ahead of the shock.

$\langle Z \rangle$ changes from 4 to 5 across the shock. The thermal equilibration between ions and electrons happens within a short spatial length in the shocked region, $\Delta x < 20 \mu\text{m}$, while no preheat is observed in the upstream region ahead of the shock. This result is in agreement with what was stated above, further indicating that electron thermal conduction alone cannot be responsible for the observed shock precursors in the experiments.

If we consider, instead, the case with the radiation transport switched on shown 5.24(b), it is noticeable that radiative losses can affect the overall shock structure. The plasma compression across the shock is stronger with ions and electrons reaching a lower maximum temperature at the shock front. As can be seen from the figure, radiation transport also impacts the plasma properties of the post-shocked region. A discussion of the stagnation layer goes beyond the purpose of this analysis, therefore we will comment only on the differences in the upstream region ahead of the shock.

A preheat region is seen in the unshocked plasma, reaching up to 12 eV near the shock transition while slowly decaying over a long distance of a few hundred microns. This increase in temperature was not observed in the absence of radiative effects [Fig. 5.24(a)] and can be attributed to the emitted

photons by the hot electrons from the shocked region absorbed by the cooler plasma inflow. The strength of this radiative precursor is significantly influenced by the opacity of the plasma, which ultimately can affect both the precursor temperature and its spatial extent. Despite the noticeable temperature increase, this preheating is not enough to further ionise the plasma inflow as, instead, would be expected during the formation of the electron density features observed in Fig. 5.22. This discrepancy with the experimental data can be better seen from the electron density profiles represented in Fig. 5.25. The figure compares the expected electron density in the pre-shock regions from the two

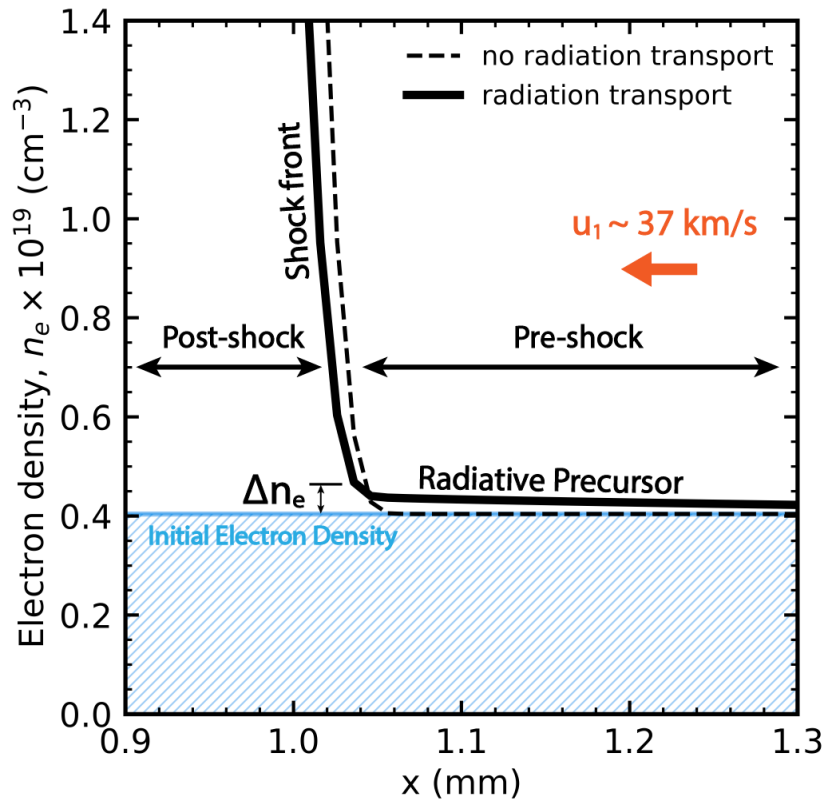


Figure 5.25: Electron density profiles of the reverse shocks presented in Fig. 5.24. A small increase in electron density is noticeable between the two cases, with and without radiation transport. The effect of radiative losses on the upstream predicted by the simulations is much lower than the experimental observations, possibly due to differences in the plasma opacity between the code and the experiments.

simulations. A small increase in electron density is observed when radiation transport is considered (solid line), however, this is not as significant as the features observed from the interferogram. This is somewhat expected as the 2nd ionisation threshold for silicon is approximately $\simeq 16$ eV which is higher than the maximum temperature of the precursor predicted by the simulation. This inconsistency is likely justified by differences in the plasma opacity between simulations and experiments rather than a misinterpretation of the radiative nature of these observed shock precursors, indicating that more

accurate simulations with various opacity models may be required.

At this stage, final statements on the observed shock precursors cannot be made without further investigations, however, these observations raise even more questions on the role of radiative cooling effects on the shock structure which will be presented in the following section. It is important to highlight that

The role of radiative cooling effects in the shocked layer

In the presented radiation transport simulations, a measure of dimensionless distance known as the optical depth, τ_ν , was assumed. For a given emitted radiation frequency, ν , the optical depth is the natural logarithm of the ratio of the incident to transmitted radiation power through the material. Large values of optical depth, $\tau_\nu \gg 1$, suggest that plasma opacity and radiation transport are dynamically important. Under this regime, the plasma is said to be "optically thick" and the emitted radiation can be re-absorbed by the surrounding plasma materials. This assumption might not necessarily hold at any spatial scales or moments during the system's evolution, resulting in strong radiative effects that can play a significant role in removing energy from the system. In this regime, the radiation power falls exponentially with optical depth and when optical depth approaches zero, $\tau_\nu \ll 1$, a plasma is said to be "optically thin" for the specified radiation, thus the opacity can be neglected. Although it is possible to experimentally estimate the optical depth of a plasma, those assessments are usually not trivial, relying on plasma spectroscopy techniques which at present are not available.

Here, a discussion on the role of radiative cooling effects is therefore presented under the hypothesis of optically thin plasma with three main questions to address: first, is radiative cooling important in this system? And if yes, how does it shape the temperature profiles immediately after the shock front? In reverse shock experiments presented in Chapter 4, accretion shocks can be affected by thermal instabilities, leading to the potential disruption of the shock system, hence ultimately developing turbulence. Are these instabilities expected to develop at the shocked layer?

The key insight emerging from comparing the data with the 1-D shock model is that an effective γ below 1.2 is required to justify the observations. As already mentioned previously, the nature of this

assumption can derive from both ionisation and cooling effects, which can strongly affect the internal pressure of the shocked layer and consequently the rate of shock accretion. As discussed in Chapter 4, the impact of radiation strongly depends on the electron temperature and density of the plasma, which are typically described by non-linear cooling curves denoted as Λ ($\text{erg s}^{-1}\text{cm}^3$). To assess the dominance of radiative cooling in the experimental observations, it is useful to compare the radiative cooling time with the characteristic time scale of the experiment. As already seen in the previous chapter, the cooling time is defined as the ratio of thermal energy to radiative power, $\tau_{cool} = U_{th}/P_{rad}$ and it can be expressed as:

$$\tau_{cool} = 2.4 \times 10^{-12} \frac{(\bar{Z} + 1)T_e}{\Lambda(n_i, T_e)n_e} \quad (5.16)$$

Here, $\bar{Z}$ represents the average charge state and T_e is the electron temperature in electron volts (eV). Assuming plasma conditions just after the post-shock region with $T_e = T_2 \simeq 20$ eV, $n_{e,2} \simeq 1.5 \times 10^{19} \text{ cm}^{-3}$, $\bar{Z} = 5$ and the corresponding value of the cooling function, $\Lambda(n_i, T_e) = 4 \times 10^{-22} \text{ erg cm}^3/\text{s}$ (Fig. 5.26), the calculated cooling time is approximately $\tau_{cool} \simeq 50$ ns. For ions to cool down significantly through radiative processes, the thermal equilibration time must be much shorter than the radiative cooling time. Additionally, the radiative cooling time should be shorter than the hydrodynamic time scale of the system to have a dynamically relevant effect. The hydrodynamic time scale, estimated as $\tau_{hydro} \simeq \Delta\xi/c_{s,2} \simeq 50$ ns, is just comparable to the cooling time and approximately one order of magnitude longer than the equilibration time.

Under these conditions, radiative cooling can potentially contribute to dissipating energy away from the shocked layer. At this point, an interesting question is whether radiative cooling is sufficiently strong to dissipate energy at the timescale of the relaxation region just behind the shock front. It is important to note that multiple processes influence the energy exchange between ions and electrons which occur at different time scales. A simple model can be used to incorporate the dynamics of thermal equilibration between ions and electrons, as well as the radiative cooling effects within the shocked layer. The governing equations for the system are as follows:

$$\frac{dT_i}{dt} = \bar{v}_e^{i|e} (T_e - T_i) \quad (5.17)$$

$$\frac{dT_i}{dt} = \bar{v}_e^{e|i} (T_i - T_e) - 10^{-7} \frac{\gamma - 1}{\gamma} \frac{\Lambda_{cool}}{e} \quad (5.18)$$

In these equations, $\bar{v}_e^{i|e}$ and $\bar{v}_e^{e|i}$ represent the ion-electron and electron-ion equilibration rates, respectively, as mentioned previously. By specifying the initial temperatures of ions and electrons in the shocked layer, T_i' and T_e' , along with the values of γ and the cooling rate $\Lambda_{cool} = \Lambda n_e$ (erg s⁻¹), the time evolution of the system can be uniquely determined. We can use the silicon cooling function

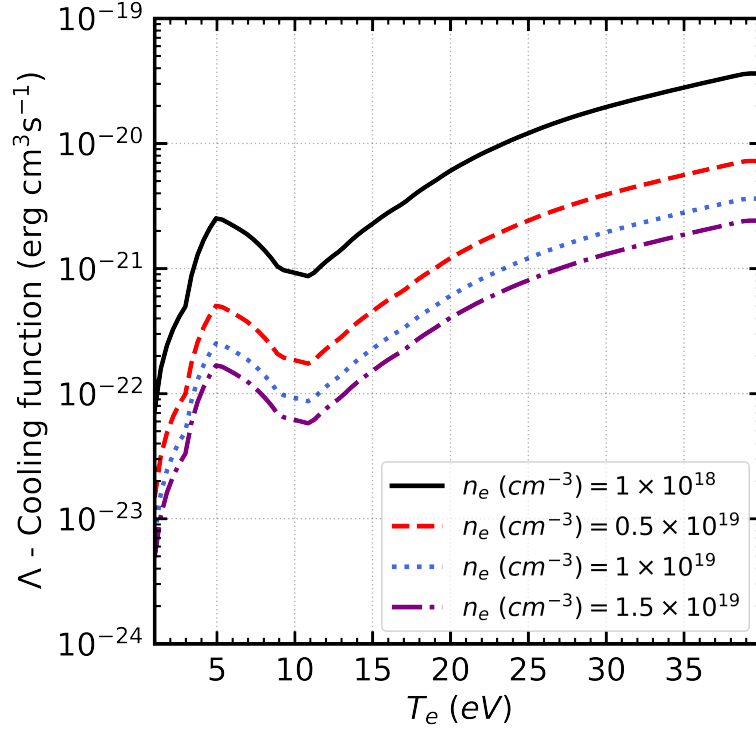


Figure 5.26: Silicon power loss curves for various electron densities as a function of electron temperature.

generated from the FLYCHK atomic code to quantify the power emitted through radiative cooling. These specific radiative power loss curves for silicon plasma are illustrated in Figure 5.26 for various electron densities. Assuming that after the shock transition, electrons are adiabatically compressed with negligible increase in temperature, $\Delta T_e \simeq 0$ eV, while ions are instantaneously shock heated, $\Delta T_i \simeq (\gamma - 1)E_i$, we can estimate the initial temperatures at the shock front as follows:

$$\begin{aligned} T_i' &= \Delta T_i + T_1 \\ T_e' &= \Delta T_e + T_1 \end{aligned} \tag{5.19}$$

The solutions of the dynamic system are illustrated in Fig. 5.27 in terms of spatial temperature profiles, assuming a constant average post-shock flow velocity of 7 km/s. In this representation, the distance origin is set at the shock transition, and larger distances correspond to deeper positions within the

shocked layer. The shaded blue and red lines represent the electron and ion temperatures, respectively, for different γ , while the data points represent the measured temperatures in the layer obtained through Thomson scattering. The measurements have been sorted based on the relative distance from the mid-plane, assuming the layer to be symmetric. The initial shock front temperatures for the considered values of γ are summarised in the following table: In the case of $\gamma = 1.2$, only half of the kinetic energy

γ	T_i' (eV)	T_e' (eV)
1.2	38	12
1.1	25	12

Table 5.4: Initial temperatures at the shock front for two different values of γ .

of the incoming flow is converted into heat after the shock, resulting in a temperature partition between ions and electrons of 40% and 10%, of the total available energy from the plasma flows, respectively. As the system evolves from its initial state, electrons quickly heat up due to thermal equilibration with ions, which occurs at a very short distance (approximately $\leq 30\mu\text{m}$) determined by the high value of $\bar{v}_e^{i|e}$. Concurrently, radiative cooling becomes significant, causing the electron temperature to reach a maximum of approximately 32 eV before steadily declining along with the ion temperature. At a distance of approximately $600\mu\text{m}$ from the shock into the layer, the ion and electron temperatures reach a value of 20 eV.

Similar dynamics occur for $\gamma = 1.1$ but with two main differences. Firstly, the energy budget converted from the kinetic energy is smaller, approximately 30% of the total available energy from the plasma flows. Secondly, the relaxation between ions and electrons occurs at a much faster rate. The temperature difference between the two particle species tends to become zero at a shorter distance into the layer compared to the case with $\gamma = 1.2$. When comparing the measured temperatures presented in section 5, it is evident that the case with $\gamma = 1.1$ agrees better with the experimental observations.

The results of the analysis suggest that radiative cooling effects can contribute to the low temperatures measured within the shocked layer on timescales compatible with the characteristic hydrodynamic scale. It should be noted that the analysis assumes no ionization processes occur within the shocked layer. However, it is challenging to separate the contributions of ionization and cooling to the estimated low γ . The fitting of Thomson scattering data indicates an average ionization $\bar{Z}$ of 5 compared to the recorded value of 4 in the upstream layer. This suggests that a significant portion of the thermal energy

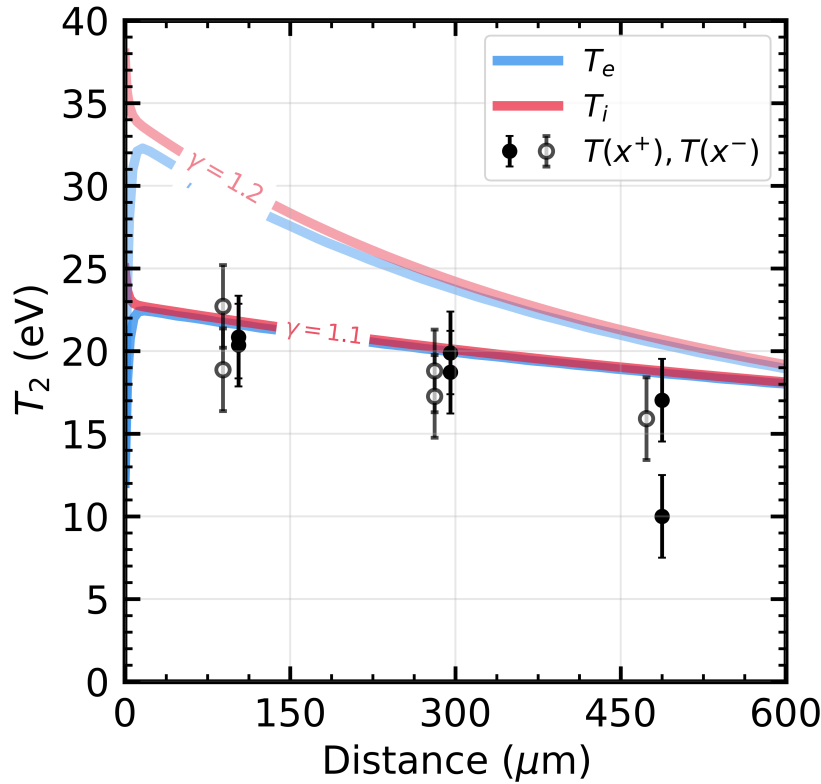


Figure 5.27: Temperature equilibration dynamics between ions and electrons in the presence of radiative cooling (solid lines) for three values of γ . The dashed lines represent the temperatures at equilibrium in the absence of any cooling effects. The initial temperature states, Table (5.4) are shown for the three considered hypotheses of γ . The electrons are adiabatically compressed, while the ions are instantaneously heated due to rapid compression.

converted by the shock may be going into ionization rather than being radiatively lost. The similarity between the hydrodynamic scale and cooling time indicates that while radiative cooling effects are present, their contribution may not be significant.

Radiative cooling instabilities

As discussed extensively in the previous chapter, radiative cooling instabilities can occur in shock systems where strong radiative losses dominate. The localized cooling leads to a decrease in plasma temperature and, consequently, a decrease in plasma thermal pressure. This compression of the cooler volume results in an increase in density and, subsequently, the cooling rate. This process can become a runaway, but it is stabilized at small length scales by thermal conduction and at large length scales by hydrodynamic motion. By performing transport calculations similar to those presented in Chapter [chapter number], we can determine the upper and lower limits of the spatial scales of perturbations

that can grow. These limits are defined by the isobaric length λ_{iso} and the λ_{field} , respectively, where the former is set by the cooling time and the latter by the thermal diffusivity of the plasma. For plasma conditions in the post-shock region with $T_i \simeq T_e \simeq 20$ eV and $n_e \simeq 1.5 \times 10^{19} \text{ cm}^{-3}$, the results are summarized in Table 5.5.

Quantity, Symbol (Units)	Si
Cooling function, Λ ($10^{-22} \text{ erg cm}^3/\text{s}$)	4
Radiated power density, P_{rad} ($10^{13} \text{ erg/cm}^3\text{s}$)	1710
Thermal energy density, U_{th} (10^8 erg/cm^3)	7.8
Cooling time, τ_{cool} (ns)	50
Electron magnetisation, $\omega_{ce}\tau_e$	0.1
Dimensionless thermal conductivity, $\kappa_{\parallel}^c[\kappa_{\perp}^c]$	7.68 [5.61]
Thermal diffusivity, $\chi_{\parallel}[\chi_{\perp}]$ ($10^3 \text{ cm}^2/\text{s}$)	9.4 [6.8]
Field's length, $\lambda_{\text{Field},\parallel}[\lambda_{\text{Field},\perp}]$ (μm)	216 [185]
Isobaric length, λ_{iso} (μm)	1040

Table 5.5: Parameters for estimating the relevant time and length scales of the cooling instabilities in the silicon plasma. The relevant plasma parameters are taken from Table 5.1. The cooling functions are taken from FLYCHK,[52] using $n_e = 10^{19} \text{ cm}^{-3}$ and $T_e = 20$ eV. The dimensionless thermal conductivities ($\kappa_{\parallel}^c, \kappa_{\perp}^c$) are obtained using the tabulated coefficients in [105].

In this case, the magnetic field is too weak to significantly affect thermal conduction, so there are no significant differences between the parallel and perpendicular directions relative to the field lines. Table [table number] provides the values of λ_{iso} and λ_{field} for the given plasma conditions. These values represent the characteristic length scales at which instabilities may develop. Understanding the growth and behaviour of these instabilities is crucial for studying the stability of the stagnation layer and the overall dynamics of the shocked system. The calculations indicate that potentially unstable modes would have spatial scales ranging between $\lambda_{\text{field}} \simeq 200 \mu\text{m}$ up to $\lambda_{\text{iso}} \simeq 1 \text{ mm}$, which are well within the observational capabilities of current laser diagnostics. However, in the experimental observations presented in this chapter, no density perturbations have been detected within the shocked layer. This raises the question of whether these instabilities are expected to grow within the timescale of the experiment, according to the thermal instability criteria. One possible answer to this question lies in the stability criteria for isobaric modes, which can be expressed as:

$$\frac{d(\log \Lambda_{\text{cool}})}{d(\log T)} < 2 \quad (5.20)$$

By applying this criterion to the cooling curves shown in Fig. 5.26, the stability conditions can be determined. Fig. 5.28 illustrates the stability analysis for the given plasma conditions, with the horizontal lines representing the thresholds for the onset of thermal instabilities. The range of temperatures where the curve is above a threshold indicates areas of stability, while instability can grow when the curve falls below a given instability threshold. According to the stability analysis, the

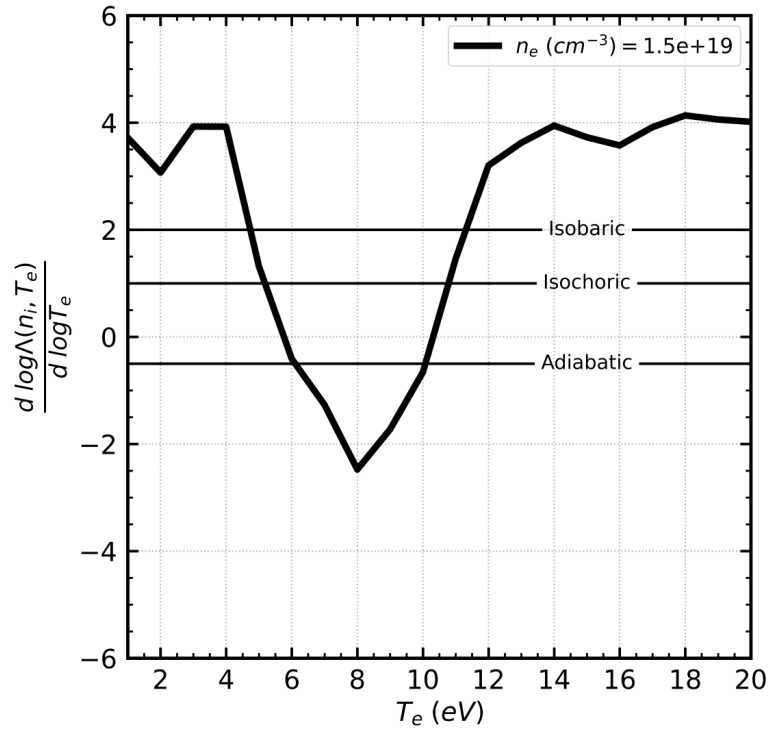


Figure 5.28: Stability analysis of the cooling function shown in Fig. 5.26 for $n_e \simeq 1.5 \times 10^{19} \text{ cm}^{-3}$. The regions below each horizontal line represent unstable regimes for the onset of isobaric (< 2), isochoric (< 1), and adiabatic (< -0.5) thermal instabilities.

system is thermally unstable in the temperature range of 5 eV to 11 eV, while it remains stable for higher temperatures. This finding is consistent with the absence of perturbations within the shocked layer, as the recorded temperatures for both ions and electrons are approximately 20 eV.

It is important to note that uncertainties in the cooling curves, particularly at such low temperatures, prevent us from drawing reliable conclusions or making predictions about the possibility of driving the system unstable under radiative cooling instability. Further experiments are necessary to explore this aspect in more detail. In summary, while the theoretical analysis suggests the potential for radiative cooling instabilities within the shocked layer, the absence of observed perturbations and the stability analysis based on the cooling curves highlight the complexity and uncertainties involved. Future

experiments and better-resolved cooling curves at low temperatures are required to gain deeper insights into the possibility of driving radiative cooling instabilities in these colliding flow systems.

5.2 Colliding counter-streaming flows in a transverse B-field

In the previous sections, the focus was on investigating the formation of a shocked layer bounded by two reverse shocks through the collision of counter-streaming flows in the presence of an ambient magnetic field. The magnetic field lines were aligned parallel to the plasma flow directions and thus perpendicular to the boundaries of the shocked layer. It was observed that the presence of this ambient magnetic field had no discernible impact on the overall shocked layer, instead only a partial contribution to the collimation of the flows is expected, since $\beta_{th} \simeq 1$ in the inflow plasma. Initial experiments conducted in [67] on a single flow propagating into a transverse ambient field have shown measurable effects of the magnetic field at the edge of the plasma outflow, changing the steepness of the low electron density profile across the plasma/vacuum interface ([67], Fig. 10).

In astrophysical accretion shocks, magnetic pressure can significantly influence the lateral expansion of the hot post-shock plasma. The compressed and heated core plasma interacting with the surrounding transverse magnetic field can cause distortion and redirection of plasma material towards the central accreting column. In the case of a strong field, ejected shells are formed and can have the potential to disrupt the reverse shock and impact the accretion processes in young stellar systems. Simulations presented in [26], show that the transverse outflows were affected by the magnetic field strength for relatively high thermal beta. This behaviour is not observed in the colliding counter-streaming flow experiments, raising questions on whether the generated ambient magnetic field can have an impact on the motion of the plasma flows at all. Possible explanations can be related to the high pressure of the plasma in the post-shock region, the estimated $\beta_{th} \simeq 5$, or the very fast field diffusion happening at the spatial scale of the interaction region due to highly resistive plasmas.

To test the hypothesis that those plasma flows can be affected by the magnetic field (e.g. resistive diffusion is not too high), an additional shot was conducted with the same experimental setup presented in § 5.1, but with the target holder rotated by 90° such that the magnetic field is now oriented

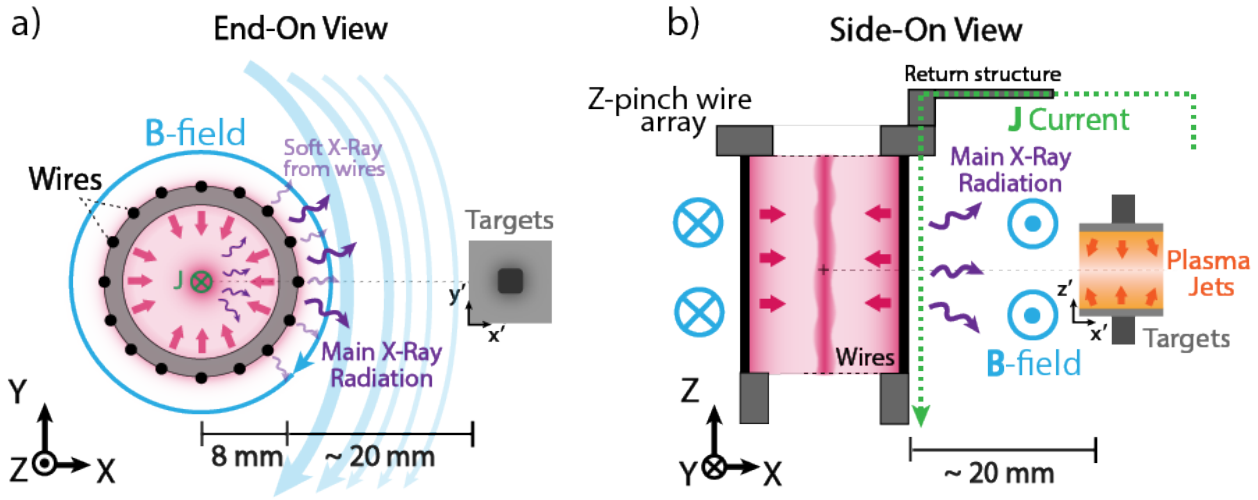


Figure 5.29: Illustration of the experimental setup of colliding silicon plasmas in an ambient magnetic field transverse to the flow propagation.

perpendicularly to the propagation of the two flows, Fig. 5.29. This configuration represents the extreme scenario, where the effects of the magnetic field on the system are expected to be the largest. In this specific experiment, the access point for the laser diagnostics is along the y direction, parallel to the magnetic field lines and achieved using side-on interferometry. In addition to the standard diagnostics, a B-dot probe is positioned opposite to the target holder at a distance of approximately 30 mm from the centre of the Z-pinch. The B-field is azimuthal, therefore it is expected that the field strength is only a function of the inverse of the radial distance from the array, $\propto 1/d$. The purpose of the B-dot probe is to measure the intensity of the magnetic field generated by the experiment at a radial position similar to the targets producing the flows while a comprehensive characterization of the magnetic field is left for future work. In this particular experiment, the x-ray pulse initiated approximately 278 ns after the current pulse.

5.2.1 Experimental Results

Figure 5.30 presents two examples of optical self-emission fast frames obtained during the experiment with silicon plasma flows in the presence of a perpendicular ambient magnetic field. In this specific shot, the first frame was recorded 132 ns (410 ns) after the X-ray pulse. The z -axis origin is at mid-target distance while the x -axis is centred relative to the visible profiles of the silicon targets. The wire array is outside the field of view with X-ray emission illuminating the solid targets from the left.

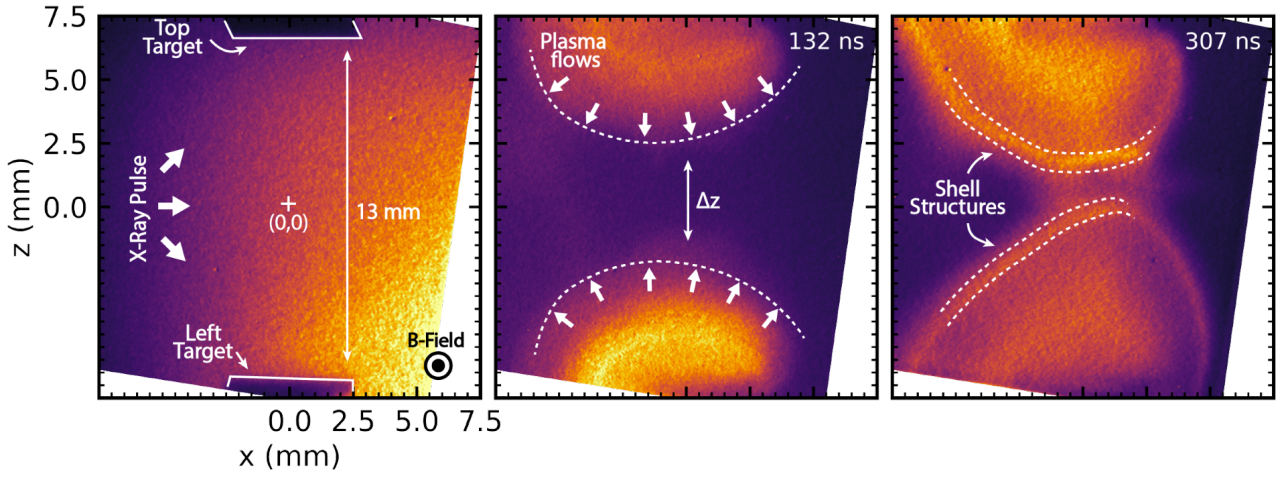


Figure 5.30: Three examples of optical self-emission frames, a pre-shot background and two images taken at 132 ns (410 ns) and 307 ns (585 ns) after the x-ray pulse, respectively. Twelve fast framing images were taken with 25 ns inter-frame starting from 132 ns.

In the first frame of the shot, the two plasma flows are already visible, with their profiles extending symmetrically up to 2.5 mm into the camera field of view. Both flow profiles are smooth and the initial flow expansion from the targets is similar to observations in the parallel ambient field experiments (§ 5.1) and in the single flow with perpendicular B-field [67]. The two plasma flows counter-propagate from the targets at similar flow speeds, as can be seen from the symmetry of the images, indicating identical outflow ablations from the target surfaces.

Approximately 50 ns after the first frame, the self-emission images reveal an intensity enhancement at the leading edges of the plasma flows. These two narrow shell-like regions, situated just ahead of the plasma flows, become increasingly distinguishable as the system evolves over time, maintaining their distinct structure until the last available frame at 407 ns (685 ns). Most noticeably, the interaction between the two flow streams does not give rise to a well-defined stagnation layer. Instead, a significant portion of the ablated plasma material remains confined within the two plasma fronts. If we compare directly with the experiments conducted in the parallel ambient field (see Fig. 5.6), the shocked layer appears 100 ns after the X-ray pulse while in this case, an early interaction is observed much later in time, $t > 200$ ns, despite having the same targets separation. Both the absence of a shocked layer and a localised increase in self-emission at plasma edges can be interpreted as evidence of a strong effect of the perpendicular magnetic field on the plasma flows. Insights on these density enhancements can be obtained by imaging the system with a laser imaging interferometry at approximately $\simeq 200$ ns ($\simeq 480$

ns) after the X-ray pulse.

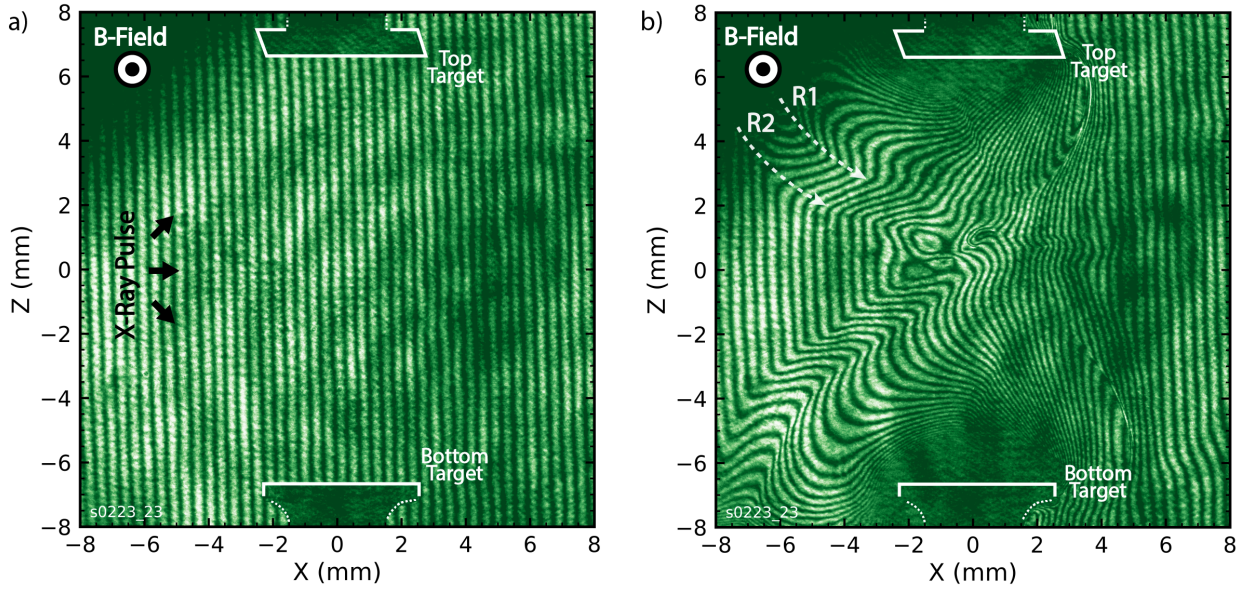


Figure 5.31: (a) Background raw interferogram taken from the side-on view. The silicon targets are at the top and bottom of the image, while the array is located on the left, outside of the field of view. (b) The raw shot interferogram taken at 200 ns (480 ns) after the X-ray pulse.

Fig. 5.31 displays the raw background and shot interferograms obtained from a 532 nm probing laser at 216 ns after the X-ray pulse (494 ns after the start of the current). The two targets are 13 mm apart from each other and about 30 mm away from the wire array located on the left-hand side of the images. In Fig. 5.31(a), the background interferogram exhibits evenly spaced straight fringes that are perpendicular to the target surface and parallel to the flow direction. This choice of fringe orientation maximizes the diagnostic sensitivity to density variations along the x-axis, where significant features are expected to manifest. Once again, the spacing between the shifted fringes is carefully chosen to achieve an adequate spatial resolution, represented by the fringe width of approximately $200\ \mu\text{m}$.

In the raw interferogram shown in Fig. 5.31(b), fringe displacement can be observed when probing through the plasma. The areas with shadows in the image correspond to regions where the probing laser is either deflected beyond the maximum acceptance angle of the optical system or absorbed. The raw interferograms reveal the presence of two distinct steep inflection points on the upper part of the plasma jets, denoted as R1 and R2. These inflections in the fringes indicate changes in the direction of the density gradients within the plasma at R1 and at the density/vacuum interface. At the time the interferograms were captured, the shell-like structures ahead of the plasma flows had already formed,

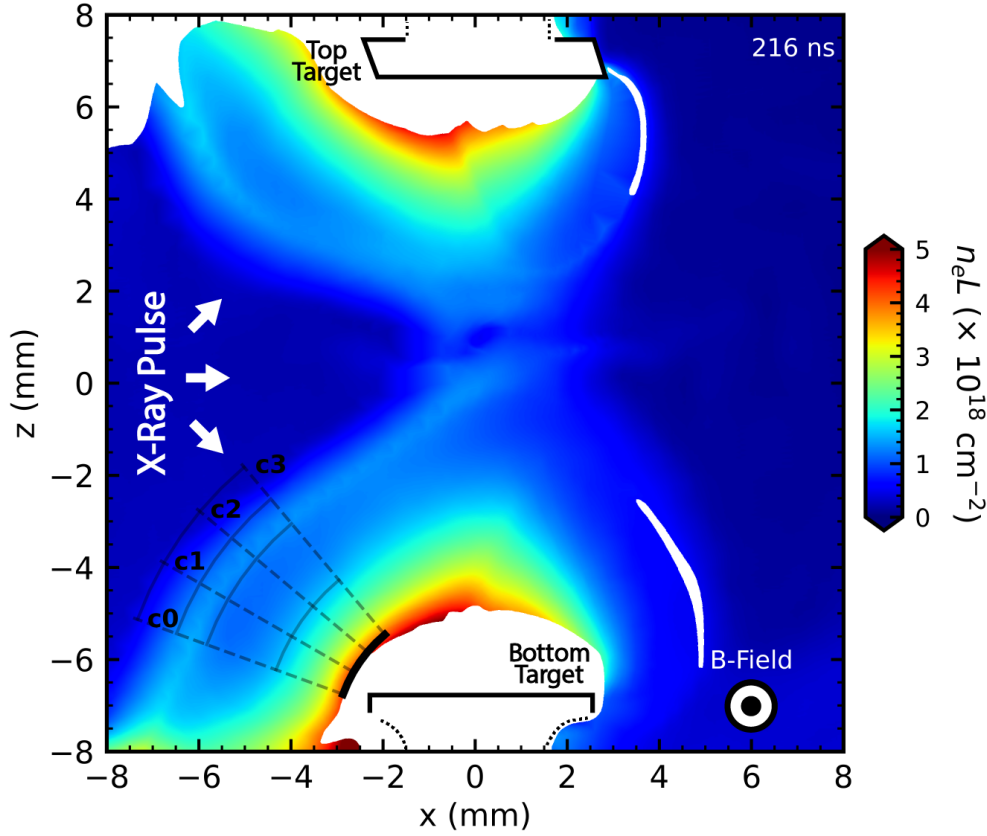


Figure 5.32: Line-integrated electron density map reconstructed from the raw interferograms presented in Fig. 5.31. The plasma flows propagate along the z -direction, transverse to the magnetic field along the y -axis. The white regions represent areas of missing data due to the loss of fringes.

as evidenced by the sequence of self-emission images.

White regions indicate areas where accurate fringe tracing was not possible. The position $x = 0$ mm corresponds to the midpoint of the target mid-plane, and the position $z = 0$ mm is set as the midpoint between the two targets. The highest electron densities are observed near the targets, with a peak density of approximately $0.8 \times 10^{19} \text{ cm}^{-3}$ at $z = -6$ mm and $z = 6$ mm. Moving away from the targets, the density of the plasma flows gradually decreases by up to 30% compared to the maximum density. Notably, at a radial distance of approximately 7 mm from both target surfaces, a dense region resembling a shock feature mediated by the magnetic field is observed. The nature of these features is challenging to determine at this stage, however for simplicity, the terms "shock" or "thin-shell" will be used interchangeably in this discussion. Fig. 5.33 shows the same line-integrated electron density measurement with a narrower heat-map scale range, such that the density enhancements are more obvious.

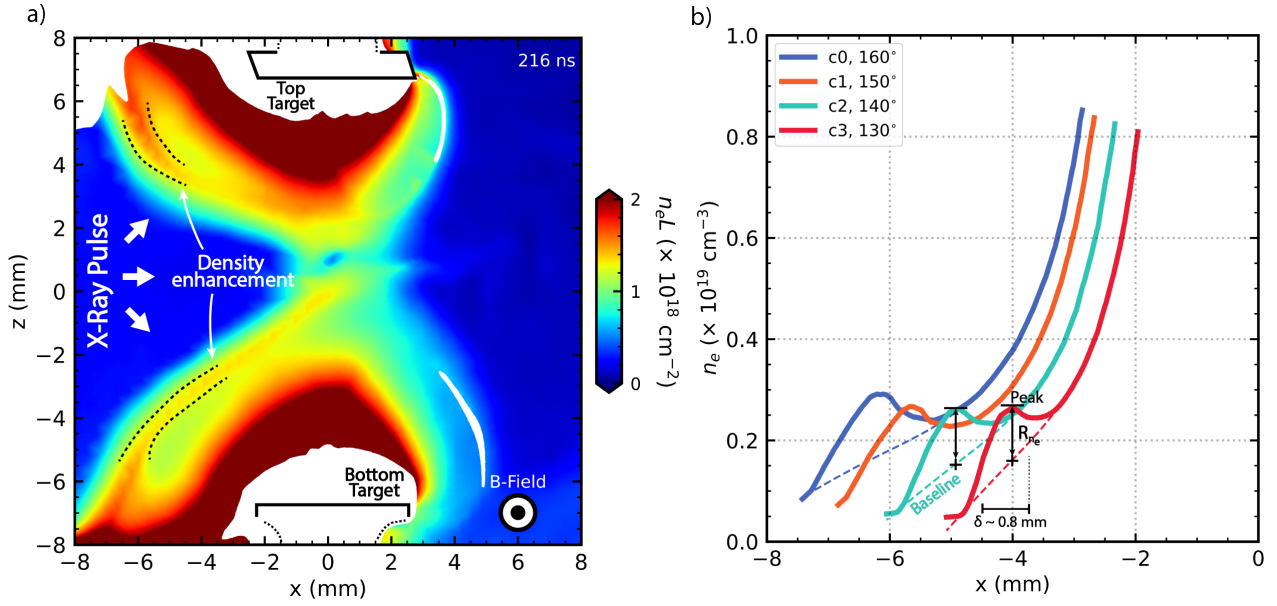


Figure 5.33: (a) Side-on Electron density map from the interferograms presented in Fig. 5.31. The heat-map scale is set to emphasise the density enhancements at the boundary of the two plasma flows. (b) Density profiles along the radial line-outs in Fig. 5.32.

The density jump observed ahead of the plasma flows can be quantitatively estimated using directly the electron density profiles shown in Fig. 5.33(b). These are taken for different radial positions from the target surface at four angles, as shown in Fig. 5.32. Different definitions of the density jump, represented by the compression ratio R_{n_e} , can yield varied estimates for shock features. Following a similar approach to what was presented in the previous sections, a local maximum electron density of around $0.27 \times 10^{19} \text{ cm}^{-3}$ is recorded while the baseline density from each line-out profile, is approximately $0.15 \times 10^{19} \text{ cm}^{-3}$, leading to a compression ratio, R , of 1.7 ± 0.1 . It is important to acknowledge that our analysis treats the inherently three-dimensional system as a two-dimensional approximation, introducing potential uncertainties in the estimates of spatial scale and compression ratio beyond the diagnostic resolution. This jump occurs smoothly over a distance of approximately 0.8 mm.

5.2.2 Possible interpretation and motivation for further work

The perpendicular orientation of the magnetic field lines relative to the flow propagation introduces new dynamics compared to the previous case with parallel magnetic fields. Notably, density features

are observed ahead of the plasma flows with the first interaction between the two flows happening at a later time than expected. Moreover, a well-defined shocked layer is not formed between the two targets, as it is, instead, observed in the case of B-field lines along the flow propagation direction. The density enhancements at the boundary of the plasma flows have not been seen in similar experiments with a single target outflow expanding into the ambient field [67], suggesting that the presence of an opposing flow has a significant effect on the ambient magnetic field and, consequently, on the system dynamics. The characteristic resistive scale of a silicon-based plasma is approximately $L_\eta \simeq 1$ mm, while the spatial scale of the plasma flows can be approximated by the target size, $L \simeq 5$ mm. At these spatial scales, the magnetic field lines do not diffuse significantly into the plasma due to its resistive scale and instead accumulate around the edges of the plasma flows.

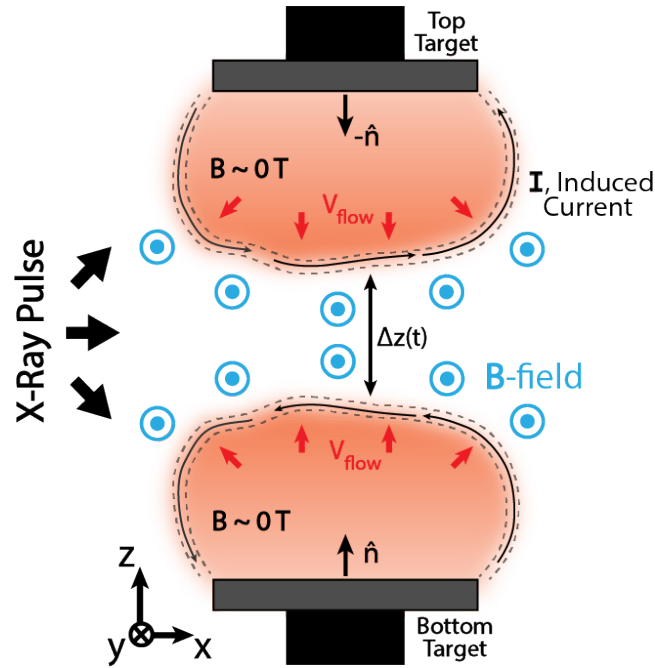


Figure 5.34: Schematic illustration of the plasma flows as they propagate in an ambient magnetic field. Two anti-parallel currents loops are formed at the boundaries of the two plasma flows, as the magnetic field is dragged towards the center. The induced currents maintain the magnetic field at almost zero within the plasmas.

As the flows approach the mid-plane, the magnetic field is swept towards the centre, causing the formation of an opposing pressure that decelerates the incoming plasma flows. The plasma from the targets is, therefore confined within a magnetic bubble surrounding the expanding flows. This magnetic bubble exhibits an increased density at the edges of the plasma flows, forming stable thin-shell structures that become denser over time. An illustration of this mechanism is shown in Figure 5.34.

Initially, the plasma flow resulting from silicon target ablation is supersonic and has higher ram pressure compared to the magnetic pressure imposed by the ambient field. This condition corresponds to a dynamic $\beta \gg 1$, allowing the plasma flows to freely expand towards the mid-plane and reach a final expansion radius of 6 mm. Over time, the magnetic pressure increases and reaches equilibrium with the ram pressure, which occurs approximately 180 ns after the start of the X-ray pulse. This equilibrium leads to the cessation of plasma flow expansion. At this stage, the pressure balance between ram pressure and magnetic pressure causes plasma accumulation at the boundary of the two flows, forming the observed shock structures. One can explain this increase in magnetic pressure at the centre of the system as field compression by the counter-propagating flows at spatial scales larger than the diffusion length. This scenario, however, is highly unlikely since magnetic field compression in a vacuum cannot occur. To compress the magnetic field a plasma medium must be present and must satisfy the frozen-in conditions, which is not the case according to the presented experimental observations. A more likely scenario, instead, can be attributed to the interaction of opposite current loops.

Plasma flows, being conductive media with high free charges, generate induced currents to maintain magnetic flux close to zero across their cross-sectional area when propagating in a static magnetic field. In our case, the plasma propagates into a perpendicular ambient field, causing induced anti-parallel current loops to close at the edge of the two plasma bubbles. Similar to two adjacent conductive wires with opposite carrying currents, these induced currents generate a repelling force at the centre of the system that results in the deceleration of the incoming flows. This can also be seen as an increase in magnetic pressure between the two plasma bubbles due to the locally induced magnetic fields. It is important to highlight that under the hypothesis of infinitely extended objects (2D system), the presence of the two current loops would have resulted in the absence of any changes in the ambient field. The geometry of the plasma bubbles, however, is naturally 3-dimensional, which may justify the increase in magnetic pressure in the central region, between the plasma flows.

At this stage, a question arises on what is the expected repelling pressure required to decelerate the flow and reach pressure equilibrium, $P_B = P_{ram}$, hence $\beta_{dyn} = 1$. The evidence that the system reaches a pressure equilibrium is supported by the sequence of self-emission images taken during the experiment. Observing the overall dynamics of the plasma flows, both plasma streams expand until 182 ns after the x-ray pulse. From that point onward, the two plasma bubbles appear stationary at a

final separation of $\Delta z(182 \text{ ns}) \simeq 1 \text{ mm}$ (measured from the origin of the x-axis). This implies that the repelling pressure must be equal to the ram pressure of the flows in the central region, otherwise, a proper hydrodynamic collision would have formed.

If we treat the repelling pressure purely as an increase in the magnetic field near the mid-plane, under the hypothesis of constant flow velocity, V_{flow} :

$$P_{ram} \approx P_{magnetic} \Rightarrow B_{max} \approx \sqrt{2\mu_0 m_i n_i} V_{flow} = 7 \text{ T} \quad (5.21)$$

where the values of the ion density and flow velocity are $V_{flow} \simeq 30 \text{ km/s}$ and $n_1 = 5 \times 10^{17} \text{ cm}^{-3}$, respectively, equal to the values measured in the previous section, 5.1. It is important to note that the assumption of constant expansion may not hold in this case, as the expansion of the two plasma bubbles slows down until the driving ram pressure is balanced by the repelling Lorentz force induced by the current loops. Nevertheless, this assumption provides a rough estimate of the magnetic field pressure required to match the experimental observations, neglecting the contribution of thermal pressure.

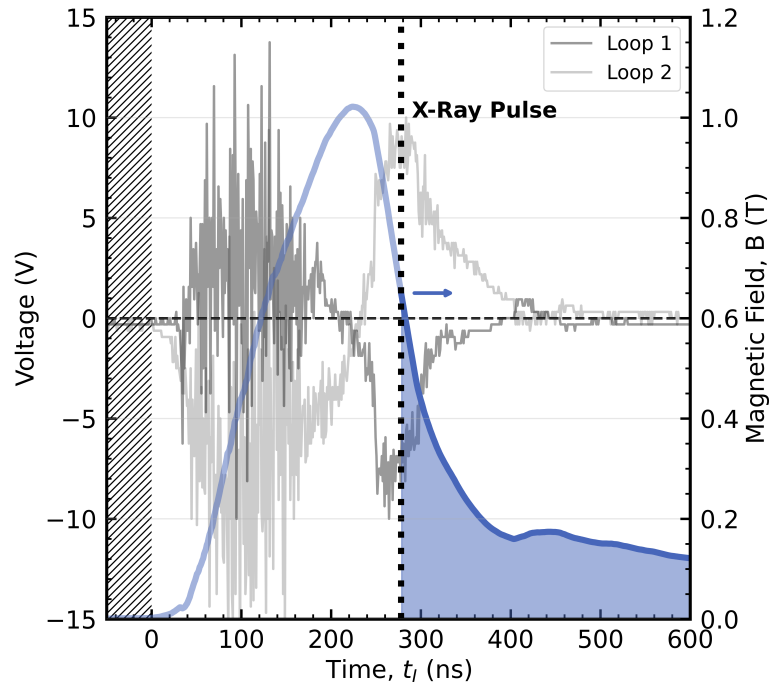


Figure 5.35: Signals from double loop B-dot probes placed 30 mm from the centre of the array. The solid blue line represented the reconstructed profile of the magnetic field from integrating the voltage signals. The magnetic field follows the current profile of the generator, peaking at 1 T field before the start of the X-pulse.

The results of the B-dot probe measurements are shown in Fig. 5.35. At the time of the start of the X-ray pulse, the measured ambient magnetic field (blue solid line) is $B = 0.7$ T, decaying rapidly with the current pulse. Based on this, the strength of the magnetic field is weak during the entire evolution of the system and more than an order of magnitude lower than what is required to affect directly the plasma flows, $B_{max} = 7$ T. For these values of ambient fields, the magnetic pressure is only dynamically important in very low-dense conditions ($< 10^{16} \text{ cm}^{-3}$) or slow plasma flows ($< 1 \text{ km/s}$), suggesting that the nature of the density enhancements is strictly tight to the dynamic of the plasma streams within the ambient magnetic field.

A comprehensive characterization of these structures cannot be achieved solely based on interferometry data and time-resolved magnetic field measurements. Additional measurements of the line-integrated magnetic field using Faraday rotation, temperatures, and velocity across the shock transition are crucial to further understand the emerging repelling pressure responsible for the flow deceleration which will be subject to future studies.

5.3 Summary of Experiments with Colliding Supersonic Flows

In this chapter, a new platform has been introduced to investigate accretion shocks in boundary-free experiments by colliding two counter-streaming supersonic plasma flows. The platform consists of the ablation by X-ray radiation from a Z-pinch wire array of two silicon targets separated by a known distance, forming two plasma outflows expanding at average speed $V_{flow} \simeq 50 \text{ km/s}$. An ambient field is naturally present in the surrounding volume, sustained by the current pulse of the generator.

In this work, a main magnetic field configuration with the field lines parallel to the direction of the flows was explored while only a few tests were conducted with field lines transverse to the plasma outflows. In the first configuration, the collision of two counter-streaming plasmas generated two reverse shocks bounding a stagnate layer, expanding backwards into the plasma upstream with field lines along the flows. The system is symmetric so that each reverse shock can be treated as an accretion shock structure with free boundaries. The properties of the shocked layer were determined by using various laser diagnostics and optical self-emission images, revealing overall consistency with a 1-D

accretion shock model for a value of $\gamma \leq 1.2$. Both ion and electron temperatures were surprisingly low, indicating that ionisation or cooling effects can likely play a role in the reverse shock dynamic. No effects of the magnetic field on the post-shocked plasma outflows suggest that the magnetic field is not strong enough to affect the dense and hot layer.

This is in contrast with many young star systems where the interaction between the core post-shock region and the lateral magnetic field can lead to the formation of dense shells due to the opposing pressure of the compressed field. Nonetheless, the connection between a typical accreting star system and the experiment can be established as long as the following underlying dimensionless criteria are satisfied:

- Sonic Mach number greater than unity, $M > 1$
- The system is collisional with ion-ion mean free path much smaller than the characteristic length, $\lambda_{ii} < L$.
- The magnetic diffusion length is smaller than the characteristic length, $L_\eta < L$
- Ram pressure of the plasma flow, $0.1 < \beta_{dyn} < 10$
- Large Reynolds number, $Re > 10^3$

According to the measured plasma parameters, all the above conditions are satisfied, representing a good comparison for accretion streams onto young stars with high beta, $\beta_{dyn} \gg 1$. A typical accreting star system has $\beta_{dyn} \simeq 1$ for a magnetic field $B \simeq 1000$ G (0.1 T). A Herbig Ae/Be star is expected to have a magnetic field of a few hundred Gauss, representing an ideal candidate for these experiments. T Tauri stars have measured fields that vary from 700 G (0.07 T) up to several kG (refs) so only the low magnetic field cases can be compared with the colliding flow experiments. To recreate in the lab the equivalent of low β_{dyn} accreting stars, it would require a magnetic field of at least 15 T at the same measured plasma flow density and velocity which is beyond the current platform capability. A side-by-side comparison of the characteristic plasma parameters between experiments and a typical accreting star system is summarised in Table 5.6.

Quantity	Symbol	Accreting Star	Solar Wind	Experiment
Flow velocity (km/s)	V_{flow}	450	$\simeq 400$	$\simeq 50$
Density (g/cm ³)	ρ	2×10^{-11}	10^{-23}	$10^{-5} - 10^{-4}$
Temperatures (eV)	T	1	12	12
Magnetic field (T)	B	0.03 – 0.3	10^{-9}	< 5
Ion sound speed (km/s)	c_s	13.6	45	$\simeq 18$
Average Ionisation	$\bar{Z}$	0.6	1	$\simeq 4$
Mach Number	M	33	$\simeq 9$	$\simeq 2$
Collisionality	λ_{ii}/L	2×10^{-3}	-	$\simeq 10^{-5}$
Magnetic diffusion ratio	L_η/L	2×10^{-7}	-	$\simeq 0.18$
Dynamic plasma beta	β_{dyn}	0.1 - 10	-	1 - 10
Magnetic Reynolds Number	Re_m	10^{10}	$\gg 1$	$5 \rightarrow 1$
Reynolds Number	Re	10^{10}	-	10^5

Table 5.6: Comparison of between solar wind and experimental plasma parameters using silicon-based targets

In addition to the configuration with a magnetic field parallel to the direction of the flows, the capability of the platform to easily change the B-field orientation allowed us to explore the case for colliding plasma flows in a transverse magnetic field. The results of the measurements using both a fast-frame camera and laser interferometry showed a very different behaviour than what was observed in a parallel configuration. The two plasma flows are significantly decelerated until density enhancements are observed at the edges of the plasma streams. These features were not observed in preliminary single-flow experiments under similar conditions, indicating that the presence of the second plasma flow has an effect on the overall system. These plasma bubbles remain stationary from the moment of the density enhancements appearance until the end of the experiment, indicating that pressure equilibrium is reached between the repelling pressure at the mid-plane and the ram pressure. If we consider the repelling force as simply an increase in magnetic field pressure, the required magnetic field to stop the two outflows, $\beta_{dyn} = 1$ is estimated to be $\simeq 7$ T, way above the measured magnetic field using B-dot probes at a similar radial position from the wire array as the targets. Moreover, no shocked layer is formed with only low dense material interacting at the mid-plane. The density jumps ahead of the plasma flows exhibit a striking similarity with MHD shock structures common in many astrophysical systems when supersonic plasmas interact with magnetospheres. Although more measurements are required to investigate the nature of these density features, the platform revealed promising results for investigating properties of plasma-magnetic field (vacuum) interfaces and solar wind interactions with magnetospheres and even in binary star systems.

Chapter 6

Conclusions and Future Work

This thesis presents investigations on reverse shocks and radiative cooling effects in high-energy-density plasma experiments relevant to astrophysical accretion shocks. Two experimental platforms both driven by the pulsed power generator MAGPIE were used to produce reverse shocks in the laboratory: a cylindrical exploding wire array with a conductive planar obstacle and the newly developed X-ray driven platform. As part of the experimental work, a new refractive-based diagnostic, named the imaging refractometer, has been also developed and applied to quantify density structures in reverse shock experiments using an exploding wire array. The imaging refractometer technique was introduced in Chapter 3 while the main experimental results of the reverse shock experiments were discussed in Chapters 4 and 5 together with their respective discussions and interpretations of the results. This chapter summarises the main conclusions and presents suggestions for future work.

6.1 Conclusions

Accretion shocks were studied in the laboratory by placing a conductive planar obstacle in a supersonic magnetised plasma flow. The plasma flow was formed from the explosion of an inverse Z-pinch wire array, producing a continuous plasma outflow lasting for hundreds of nanoseconds with an embedded magnetic field. This plasma source has been well-characterised in previous MAGPIE experiments, employing aluminium wires as the load material of the array. This thesis extended the

investigation conducted on reverse shocks in aluminium plasma to high-Z elements, such as tungsten, where radiative cooling effects were expected to be stronger. The collision of a supersonic magnetised plasma flow against a planar obstacle leads to the formation of a reverse shock, accreting as more material accumulates at the obstacle surface. This reverse shock structure appeared different between experiments with aluminium and tungsten plasma flows despite similar upstream parameters, such as mass density, flow velocity and magnetic field. In aluminium plasma, a stand-off shock was observed ahead of the obstacle, supported by the pile-up of magnetic field and the increase in thermal pressure. In contrast, experiments with tungsten plasmas revealed the absence of a stand-off shock and instead the presence of density perturbations at various spatial scales, extending as the post-shock region in aluminium plasma with characteristic length comparable to the magnetic diffusion length of the plasma ($L_\eta \simeq 3$ mm). In both experiments, a stagnation layer is also observed near the obstacle surface characterised by the stagnant dense plasma.

This striking difference in the reverse shock morphology reflects the importance of radiative effects in astrophysical accretion shocks which do not exhibit the typical one-dimensional profile, limiting our current understanding of accretion columns in young stellar objects. Characteristics of these systems, such as mass accretion rates, are typically inferred from direct optical observations by assuming uniform flow streams and plasma conditions of the post-shock regions stationary over time. These simplistic assumptions combined with the current limitations of numerical study due to lack of experimental data have been the main driver to replicate accretion shocks in highly controllable environments.

Detailed spatially and temporally resolved measurements were performed in experiments with tungsten plasmas using the MAGPIE diagnostic suite, including laser interferometry, optical Thomson scattering and optical self-emission imaging. These measurements allowed to quantify tungsten plasma parameters, such as electron density, flow velocity and an upper estimate plasma temperatures, in both free-flow conditions and in the presence of the planar obstacle which was then compared with aluminium plasma. In addition to laser-based diagnostics, physical B-dot probes were used to measure the plasma frozen-in magnetic field. This was found consistent with measured values of the magnetic field in previous aluminium experiments, indicating that the pile-up of the transverse B-field can still occur and is dynamically important to shape the reverse shock morphology.

To further investigate the nature of density perturbations observed in tungsten plasma flows, the novel imaging refractometer technique was used in parallel to a standard shadowgraphy imaging system. This new refractive-based diagnostic allowed to directly measure laser deflection angles along one direction (z-axis) when passing through a non-uniform plasma. The characteristic size of these density perturbations was estimated assuming random walk and isotropic perturbations. The largest spatial scale of the density structures was found to be $\delta_y \leq 200 \mu\text{m}$ in the direction parallel to the magnetic field, while almost three times smaller in the transverse direction of the B-field due to partially magnetised electrons. These estimates were much smaller than the magnetic diffusion length while consistent with the expected changes in electron thermal conduction along and perpendicular to the magnetic field lines.

Combining these results with measurements from other diagnostics it was found that the nature of these density perturbations is likely associated with strong radiative cooling effects in tungsten plasmas, specifically due to thermal instability. This hypothesis was also bench-marked with the stability criteria provided by Field's analysis which agreed with the observed behaviour of both aluminium and tungsten flows at the measured temperatures. Although radiative effects seem to be more important in tungsten flow than aluminium plasmas, cooling instability could still develop in regions where density is large enough and temperatures are within a specific range. This could be the case of plasma materials in the stagnation layer, where structures were observed with a growth rate slower than in tungsten plasma. It is important to note that, despite this evidence, it is not possible to draw definitive conclusions on both the threshold and growth rate of the thermal cooling instability due to a lack of accurate cooling functions. These experiments revealed insights into the complex and rich physics behind accretion shocks along with the importance of radiative cooling effects on the shock dynamics. In addition, the presence of a relatively weak transverse magnetic field had also a strong effect on the reverse shock structure, despite the high dynamic β of the plasma flow, preventing us from making direct comparisons with astrophysical systems.

To overcome the above and other limitations of traditional inverse wire arrays, a new platform based on X-ray-driven ablating targets was developed, generating boundary-free accretion shocks from the collision of counter-streaming flows in an ambient magnetic field. The X-ray radiation was produced by a Z-pinch wire array while the ambient field was naturally sustained by the current pulse of the

generator. In this thesis, the magnetic field configuration with field lines parallel was mainly explored while only a few tests were conducted with magnetic field lines transverse to the plasma outflows. In the former configuration, the collision of two counter-streaming plasma flows produced two reverse shocks bounding a stagnate layer, which expanded into the upstream with field lines along the flow. The symmetry of the system allowed to treat each reverse shock as an accretion shock structure with free boundaries, which could be then compared with the 1-D accretion shock model. Properties of the shocked layer were determined by various optical diagnostics, such as optical self-emission imaging, laser interferometry and Thomson scattering. Overall, the results of the analysis revealed consistency with the 1-D model under values of $\gamma \leq 1.2$. The estimated low value of the polytropic index was also justified by the measured low plasma temperatures suggesting that cooling effects and ionisation could be important in the reverse shock dynamic. In addition, electron density features were observed ahead of the reverse shocks, suggesting the presence of radiative precursors. The radiative nature of these shock precursors was further supported by numerical simulations using the 1-D code HELIOS-CR. Surprisingly, no effects of the magnetic field on the transverse shocked plasma outflows were observed, which indicates that the magnetic field is not strong enough to affect the dense layer. Although this is in contrast to many young star systems, these experiments are still comparable to astrophysical accretion streams since the measured plasma parameters satisfied all the dimensionless scaling conditions.

Finally, the flexibility of the platform to easily change the orientation of the B-field allowed to test the effect of a transverse ambient field on colliding plasma flows. The results from optical self-emission imaging and laser interferometry data showed significant deceleration of the two plasma flows and no formation of a shocked layer. Density enhancements at the edges of the plasma streams were observed, instead, resembling MHD shocks common in many astrophysical systems when supersonic plasmas interact with magnetospheres. Further characterisation of these density features will be part of future work. To conclude, these results successfully demonstrated the capability of the X-ray driven platform to produce plasma systems scalable to astrophysical objects, providing new avenues for the laboratory of astrophysics experiments.

6.2 Suggestions for Future Work

Three possible ideas for future work are described here, primarily leveraging the novel X-ray driven platform. Beyond these propositions, several elements of the work detailed in this thesis remain under active exploration. Notably, there is a current absence of numerical simulations on colliding plasma flows for both magnetic field configurations. These simulations could provide additional evidence for the radiative nature of the observed reverse shocks and a better understanding of the impact of the ambient magnetic field on the shocked layer.

1. In this thesis, we presented only experiments using silicon-based colliding plasma flows in both magnetic field configurations with field lines parallel and perpendicular to the plasma flow direction. Exploring the interaction of counter-streaming flows with other target materials can help to further understand the role of ionization, radiative effects and magnetic fields on the shock structure, as well as on the potential capability to produce faster plasma flows. Preliminary experiments conducted with carbon targets have shown negligible differences in the flow velocities despite the lower atomic mass, raising questions about the ablation process driving the plasma outflows. Additionally, supersonic colliding flows from mixed target materials (e.g. carbon-silicon) can be performed to investigate mixing and reverse shock structures in the case of multi-species plasmas. In this case, measurements of the plasma self-emission lines in the interaction region can also be recorded thanks to the open-access geometry of the X-ray driven platform. All these experiments can help to better understand reverse shocks under different plasma regimes and in general, characterise plasma flows generated from ablating targets.
2. A possible next step is to investigate the stability of shocked layers from colliding plasma flows. Experiments with silicon targets revealed very stable reverse shocks expanding as more material accumulates at the central region. Open questions remain on whether these layers can be disrupted by thin-shell instabilities under specific plasma conditions. The use of metallic targets as plasma sources could allow to access strong radiative cooling regimes due to the high atomic number, potentially showing the growth of unstable modes at the stagnation layer. This will require a major redesign of the platform hardware in order to minimise the interaction between

the metallic targets and the ambient magnetic field, which could introduce unwanted effects on the overall system dynamics. Alternatively, solid targets can be manufactured to present well-defined corrugations. This could result in less homogeneous leading edge profiles of the plasma flows, hence possibly driving the shocked layer unstable.

3. Finally, studying turbulent plasmas in a highly controlled environment is an active area of research within the HEDP community. While the majority of existing research has been conducted using laser facilities, the brief duration of these experiments (a few nanoseconds) limits the range of physical mechanisms that can be explored. The X-ray driven platform, however, has demonstrated its ability to drive plasma flows that are scalable to astrophysical objects over an extended period of hundreds of nanoseconds, which makes it an ideal driver for long-lasting turbulent plasma experiments. Potential avenues for achieving this include the use of corrugated targets and solid meshes or grids with varying spatial scales. Furthermore, the possibility of generating turbulent plasmas in a laboratory setting could significantly enhance the development of the imaging refractometer technique as a potential quantitative plasma diagnostic for highly non-uniform plasmas.

Appendix A

Analysis of the shock-bounded stagnation layer

A reverse shock system is formed when a supersonic plasma flow interacts with a hard surface which provides a supporting pressure for the stagnant plasma. In tungsten experiments, optical self-emission images reveal the reverse shock as a dense thin layer in the proximity of the obstacle surface expanding towards the wire array as more material accumulates at the stagnation layer. This shocked layer is bounded on one side by the shock front and on the other side by the stagnant plasma, exhibiting a hydrostatic structure that increases in thickness as the system evolves over time. A summary of the reverse shock features in tungsten plasma is shown in Fig. A.1 from an example of plasma self-emission image taken at 464 ns after the current start.

Thin layers of matter are common in many astrophysical shocks [38, 39] and in the laboratory [40, 41] whenever radiative losses of energy become significant. These layers tend to accumulate mass as they move, becoming unstable if the internal pressure acting always normal to the layer surface is not balanced with the ram pressure of the incoming flow. This can occur in the presence of a modulation in the layer, where this imbalance causes the lagging material to decelerate less rapidly than the leading material at a lower density, inverting the modulation, and producing either a power law or exponential growth. This unstable mechanism was originally theorised in [42] and confirmed experimentally in [43]. In more recent literature, this shock instability is often referred to as Vishniac-Ryu instability [34]

in the context of decelerating shocks or "thin-shell instability" and "nonlinear thin-shell instability" in colliding plasma flows [36, 44, 45] for linear and non-linear regimes, respectively.

In our experiments, the shocks do not propagate in an ambient medium as typically described in the literature, but instead, their motions are strongly coupled to the change in ram pressure due to the finite current pulse and the limited mass of the ablating wires. As less plasma outflow is brought into the stagnation layer, the shocked layer appears to decelerate upstream until reaching its final position at the late time of the experiment. Despite this subtle difference between the astrophysical cases and the current experiments, the stability criteria for shock-bounded layers should not be dissimilar between the two systems.

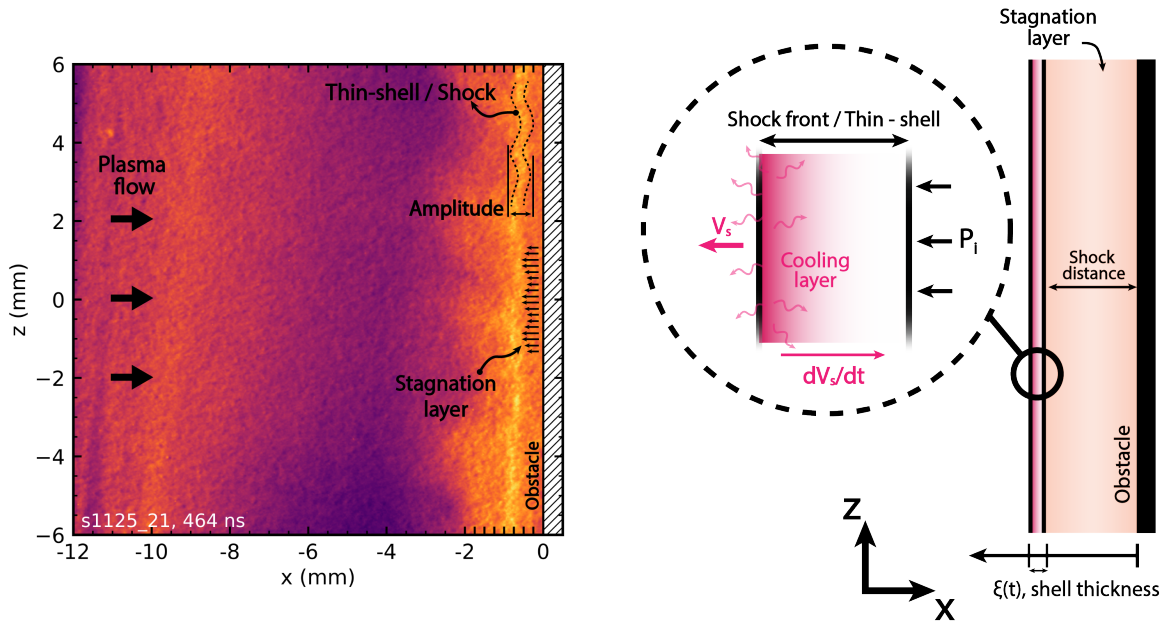


Figure A.1: (Left) Example of optical self-emission image of the xz -plane at 464 ns taken during s1125_21 shot. (Right) A 2-dimensional illustration of the shocked layer, bounded by the reverse shock front and by the supporting pressure of the stagnant plasma, P_i . Any cooling effects are confined in a very narrow region behind the shock front.

The 2-dimensional shock-bounded layer model used in this discussion is illustrated in Fig. A.2, showing both density and temperature profiles across the shocked region. If we assume the shock front at position x_0 , the thin shell layer can be defined as the post-shock region supported by the uniform pressure P_i of the stagnant plasma near the obstacle surface at its rear edge, x_i . As more plasma material is shocked at the leading edge, the layer thickness $\xi(t)$ increases over time, starting from an initial value of ξ_0 . Most of the kinetic energy from the incoming flow is carried by the ions because

of the large mass, and therefore they are the ones heated at the shock transitions by viscous effects. The incident electrons, instead, are negligibly heated, typically through adiabatic compression. At the shock transition, the heating process occurs instantaneously, with ions and electrons reaching the temperatures T'_i and T'_e , respectively. These values represent the final temperatures of the post-shock region, in the absence of any energy losses or energy exchange between electrons and ions. If energy exchange mediated by Coulomb collision and radiative losses occurs, a relaxation region is formed behind the shock front, with length Δx . Here, ion and electron temperatures equilibrium until reaching a final temperature T_2 . For those cases where these effects occur at very short time scales, $\Delta x \ll \xi$, the shocked layer can be considered isothermal, characterised by negligible temperature gradients. Plasma density also increases in shocks, from upstream density ρ_1 to immediate post-shock value ρ_2 . Under the hypothesis of isothermal layer and in the presence of deceleration, the density decreases exponentially away from the shock front, until reaching the value ρ_i .

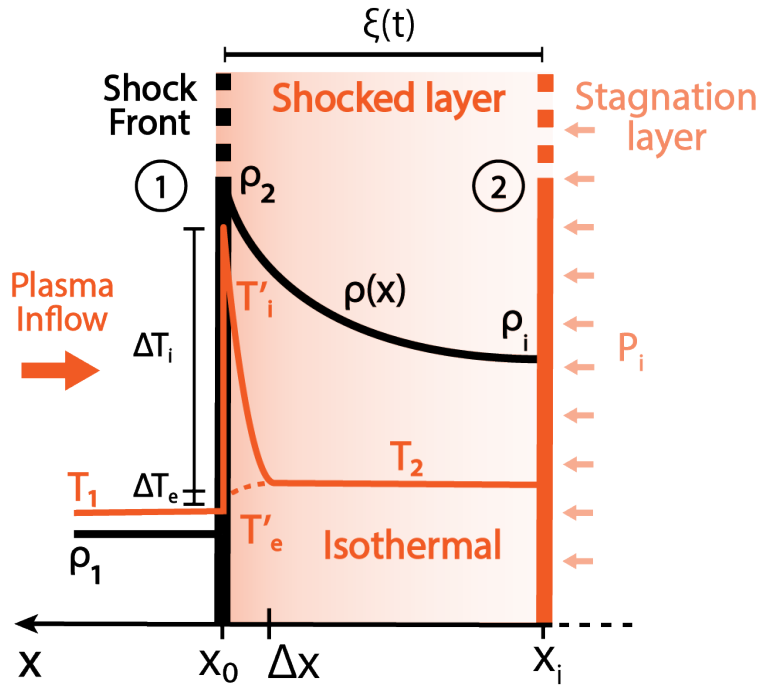


Figure A.2: Temperature and density profiles across the shock-bounded layer. Thermal equilibration and radiative cooling effects happen in a narrow region behind the shock front at spatial scales much shorter than the overall layer thickness, $\Delta x \ll \xi$. Overall, the shocked layer is isothermal at the final temperature, T_2 , and the density profile, $\rho(t)$, decreases exponentially. The supporting pressure, P_i acts on the rear edge of the thin shell, given by the accumulated stagnant plasma near the obstacle surface.

It is important to note that there are many other possible density profiles that can lead to quasi-steady hydrostatic layers, depending on the heat transport dynamics. In general, the density in the

shocked layer may be determined by physics other than simple hydrodynamics and might not be well described by any values of the polytropic index γ . For simplicity, here we will discuss only the stability of isothermal shock-bounded layers, as originally described in [34, 36].

From plasma self-emission images presented in the previous chapter (Fig. 4.12 and Fig. 4.13), the reverse shocks in tungsten experiments are strongly modulated along the entire height of the shocked layer (z-axis). This modulation starts at $\simeq 260$ ns after the current pulse and it grows in amplitude until the very last available frame ($\simeq 585$ ns). This modulation appears monochromatic with a characteristic wavelength $\lambda_0 \approx 3 \pm 0.5$ mm measured from the analysis of intensity profiles along the shocked layer discussed in the previous sections and shown in Fig. 4.16. Similar modulations were also observed along the stagnation shock in aluminium experiments with the key difference that the growth there was much slower. As shown in Chapter 4, this is not surprising as less plasma flow is able to reach the shocked layer due to the presence of the stand-off shock. Based on these considerations, natural questions arise on the growth mechanism of this observed modulation and in general on the stability of the reverse shocks according to the thin-shell instability theory.

To answer these questions, the plasma parameters of both the inflow and shocked layers must be known. The former has been measured and summarised in Table 5.1, while the values of the plasma state within the shocked layer are not available. This shock feature is not resolvable by the current laser imaging diagnostic as strong density gradients in that region preclude any direct measurements of the shock. The only available data are mainly qualitative represented by the self-emission images. Therefore, in this discussion, a few assumptions are required.

- The reverse shock is a strong shock. The plasma flow at the end of the observed deceleration region is still supersonic with a velocity of $V_1 \simeq 20$ km/s.
- The plasma within the shocked layer is relatively cold and isothermal. The ion and electron temperatures are in equilibrium, and the final temperature is assumed to be $T_2 = T_{i,2} = T_{e,2} \simeq 10$ eV. The average charge state is unchanged $\bar{Z} = 6$ (estimated using FLYCHK code) and the speed of sound in the layer is $c_{s,2} = \sqrt{k_B T_2 (1 + \bar{Z}) / m_i} \simeq 6$ km/s.
- Any radiative cooling effects and equilibration processes occur at very short timescales, and

hence are confined in a narrow region behind the shock front. The tungsten flow is strongly radiative, as will be discussed later in the section, with a cooling time of $\tau_{cool} \simeq 7$ ns and a thermal equilibration time of $\tau_{ei}^e \simeq 5$ ns. In both cases, the relaxation region is very short relative to the layer thickness, $\Delta x \simeq 30 \mu\text{m} \ll \xi$.

- The effective polytropic index γ is assumed to be small, $\gamma \leq 1.2$. This is a reasonable assumption in systems where radiative cooling effects are important. Based on this upper limit and the measured upstream Mach number, $M_u = u_1/c_s \simeq 3$, the expected density and velocity jump are estimated to be $R = \rho_2/\rho_1 = u_1/u_2 = 5$.
- The density profile across the shocked layer decreases exponentially. This is the standard case assumed in the classic description of shock-bound layer instability [34] which will be the only focus of the discussion. Other density profiles and associated stability criteria have been explored in [37].

A summary of the plasma parameters considered in this section is presented in Table A.1. It is important to note that, despite the assumptions above being reasonable according to experimental observations, these may not accurately represent the true plasma state of the shocked layer. Additional experiments and numerical simulations are required to further support these assumptions.

Quantity	Symbol	Inflow (at $x = -2$ mm)	Thin-Shell (inferred)
Flow velocity (km/s)	u_1	20 ± 1	4
Electron density ($\times 10^{18} \text{ cm}^{-3}$)	n_e	2.5	12.5
Mass density ($\times 10^{-5} \text{ gr/cm}^3$)	ρ	5	25
Temperatures (eV)	$T_i \simeq T_e$	< 10	10
Average ionization	$\bar{Z}$	6	6
Sound speed (km/s)	c_s	6	6
Polytropic index	γ	≤ 1.2	≤ 1.2

Table A.1: Plasma parameters in the upstream and in the shocked layer of tungsten reverse shock. The plasma state within the shocked layer is inferred from the measured upstream values, considering the hypothesis of strong radiative shock and isothermal layer.

The stability criteria for shock-bounded layers

In tungsten experiments, the shock-bounded layer confined on one side by the reverse shock front and on the other side by the stagnation layer is accompanied by the presence of a well-defined modulation along the layer with a characteristic wavelength of $\lambda_0 \simeq 3 \pm 0.5$ mm. According to [34, 36] the stability of thin layers depends on both the internal properties of the layer and its thickness, $\xi(t)$. Thicker layers generally exhibit greater stability against wave disturbances propagating along the layer than thinner layers. Moreover, only perturbations with wavelengths longer than the layer thickness develop, therefore it is important to evaluate the thickness of the shock-bounded layer over the duration of the experiment.

In the previous section, estimates of the layer thickness were characterised by horizontal intensity line-outs across the shock feature observed in the optical self-emission images, resulting in the thickness profile shown in Fig. 4.15 and again reported here in Fig. A.3 (Left). The measurements show a linear increase in thickness of the shock-bounded layer with an average expansion rate, $\simeq 1.6 \pm 0.5$ km/s, similar to the reverse shock velocity, $v_s = 1.7 \pm 0.5$ km/s [Fig. 4.14(c)]. This behaviour aligns well with the expected thickness predicted by the strong shock theory and derived in chapter 1, $\Delta\xi \simeq c_{s,2}^2 \Delta t / V_i$, assuming the values presented in A.1 for the post-shock speed of sound and inflow velocity. As the thickness of the layer evolves over time, the set of perturbations that can grow during the timescale of the experiment also changes, with only longer wavelengths, $\lambda > 2\xi$, being enhanced, while shorter modes are generally damped by layer tension. Furthermore, shorter wavelengths exhibit faster growth compared to longer wavelengths, with the upper limit determined by the speed of sound $c_{s,2}$ [44].

Under these conditions, a map of the possible unstable modes for the observed shock-bounded layer can be derived as a function of the experiment time, as shown in Fig. A.3(Right). The green region indicates modes that are stable for the given profile of the layer thickness. The modes are either damped over time or are not strong enough to disrupt the layer within the time scale of the shot. The region of possible unstable mode is defined by the excluded area in the plot, confined between the two stability limits, $\lambda > 2\xi$ and $\lambda < c_{s,2}t$. This narrow region expands over time, revealing that perturbations of a few millimetres can evolve and become unstable during the time frame of the experiment. From plasma self-emission data, the observed mode with $\lambda_0 \simeq 3 \pm 0.5$ mm begins to

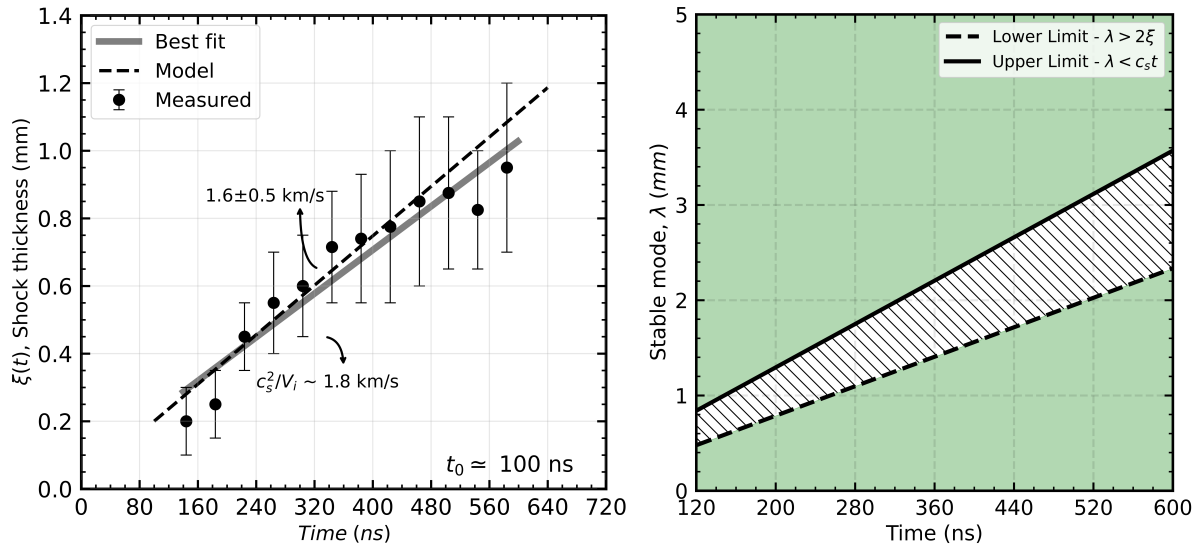


Figure A.3: (Left) Time-resolved layer thickness from optical self-emission images (s1125_21) as presented in Fig. 4.15. The dashed line represents the expected layer thickness in the strong shock limit, where V_i is the measured upstream flow velocity and c_s is the speed of sound within the post-shock region. (Right) Stability map of shock-bounded layers due to thin-shell instability. The green regions represent stable modes with characteristic wavelength λ . The excluded area represents waves that are allowed to grow along the shocked layer with thickness ξ and speed of sound $c_{s,2}$.

develop at approximately 260 ns, which is inconsistent with what is expected from the theory as only modes with $\lambda = 1 \div 1.5$ mm can grow at that time. According to the stability criteria, the measured perturbation falls within the range of unstable wavelengths at later times, specifically for $t > 400$ ns. It is worth noting that although waves develop along the shock profile and the thin-shell layer appears modulated, the system does not reach the point of instability. The stagnation shock remains intact throughout the duration of the experiment. This implies that the timescale at which the growing modes could disrupt the system is very long, or in other words, the conditions of the thin-shell layer are such that the system is actually stable at these perturbation scales. To further characterise the observed wave, it is necessary to estimate how the wave amplitude changes over time and compare it with the expectations of the theory presented in Chapter 1.

Modulation amplitude and growth rate in the linear regime

In tungsten reverse shock experiments, the visible modulation grows over time as more material reaches the shock-bounded layer. This increase in amplitude, however, never disrupts the layer over the timescale of the experiments, indicating that the growth rate of this perturbation is slow or

the mechanism driving this growth is suppressed by the gradual reduction of plasma ram pressure from the wire array. Thin-shell instabilities in a shock-bounded layer are driven by the imbalance between the ram pressure of the incoming flow and the internal pressure of the shocked layer which results in a power-law or exponential growth of unstable modes. Therefore, insights into the physical mechanism responsible for the observed perturbation growth can be obtained by comparing the increase in amplitude with the expected growth profile in the case of unstable thin-shell layers.

According to [34, 36] and more recently [37], the density profile within the shock-bounded layer plays a key role in layer stability as this is determined by the delicate balance between internal and external pressures. It is intuitive that thin layers can be susceptible to perturbations based on their "stiffness". A more "rigid" layer is less prone to perturbations and thus more stable than less stiff layers. This property is closely related to both layer thickness and internal density distribution which can depend on the heat transport mechanisms in the layer formation. Under the assumptions presented earlier of isothermal slabs, the density profile across the layer follows an exponential law given by $\rho(x) = \rho_2 e^{-x/h}$, where ρ_2 is the post-shock density and h is the isothermal scale height, $h = c_{s,2}^2/a$, inversely proportional to the layer deceleration $\mathbf{a}$. As a result, the density profile within the layer decreases along x from the maximum post-shock density near the shock front, ρ_2 , until the rear edge of the shell, ρ_i , as shown in the schematic of Fig. A.2. Taking into account the change in layer thickness over time, $\xi(t) = \xi_0 + c_{s,2}^2 V_1^{-1} t$, this leads to an expression that describes the change in density of the layer over time, as more material is shocked during the experiment:

$$\rho(t) = \rho_2 e^{-\xi(t)/h} \quad (\text{A.1})$$

The result is presented in Fig. A.4, assuming a constant deceleration of the layer, $a = 10^7 \text{ km/s}^2$, as derived previously from the reverse shock analysis, Fig. 4.14(c). The initial value, ρ_2 is determined assuming a density jump, $R = \rho_2/\rho_1 = 5$, based on the choice of $\gamma \leq 1.2$ and the measured upstream sonic Mach number, $M_u \simeq 3$. For $t > 500 \text{ ns}$, the assumption of constant decelerating motion is not any more valid as the layer reaches its final position from the obstacle. At this point, it is assumed that the layer density remains constant at $\simeq 5 \times 10^{-4} \text{ gr/cm}^3$.

To compare the observed amplitude of the modulation with the expected growth rate from the

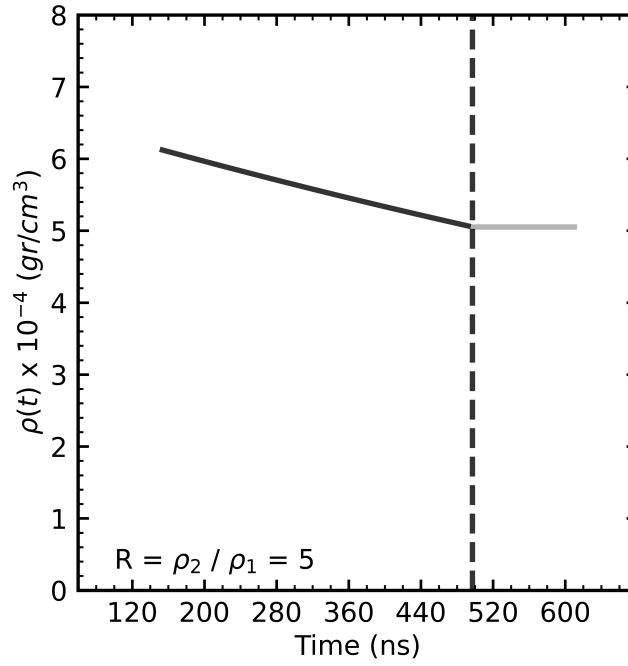


Figure A.4: Change of the internal layer density over time as more material is shocked, contributing to the increase of the layer thickness. This density profile is obtained assuming a 5-fold compression at the shock transition and an exponentially decreasing density distribution across the layer.

theory, the simple thin-shell instability model for shock-bounded layers confined by two accretion shocks can be applied. For details on the model see chapter 1. Although the original model was derived for two counter-streaming plasma flows, it can be modified for a single plasma stream by assuming that the dominant of one inflow pressure. This results in the drop of factor 2 in the ram pressure contribution:

$$A(t) = \frac{\lambda_0}{\pi} \sqrt{\frac{n_{i,1} V_1 (t - t_{0,\lambda_0}) + \rho_0 \xi_0 - \rho(t) \xi(t)}{\rho(t) \xi(t)}} \quad (\text{A.2})$$

This equation relates the amplitude growth of a monochromatic perturbation, denoted as $A(t)$, to the properties of the thin-shell layer. By knowing the average thickness of the layer, represented by $\xi(t)$, and the ion mass density of the slab, $\rho(t)$, we can compare the predicted growth with the experimental observations in the linear instability regime. In this context, the amplitude of the modulation is defined as the peak-to-peak spatial difference as shown in Fig. A.1(Left). Measurements of the amplitude have been performed by visually tracking the maximum and minimum peak of the perturbation in each available frame. The results are shown in Fig. A.5. The smallest amplitude of approximately $300 \mu\text{m}$ is recorded at around 260 ns. The perturbation then grows over time, reaching a maximum of 1.2 mm towards the end of the experiment. It is important to note that both low-resolution and line-integrated

effects affect strongly the accuracy of the measurements, resulting in large error bars. The amplitude model expected by the thin-shell instability theory is shown as the solid blue line. Overall, the model is consistent with the measured amplitude growth profile of the perturbation with a wavelength of $\lambda_0 \simeq 3 \pm 0.5$ mm observed in the experiments for the given set of plasma parameters summarised in Table A.1.

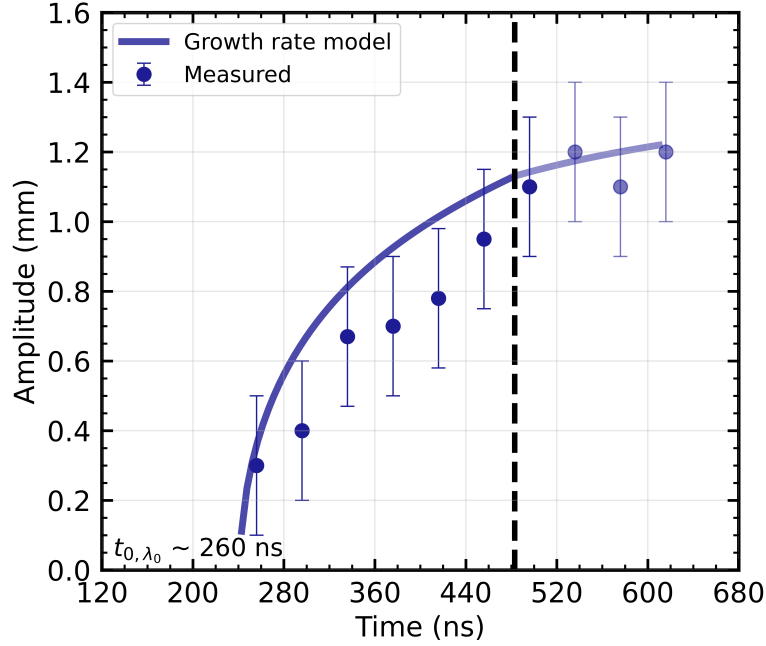


Figure A.5: Amplitude growth profile of the observed perturbation with $\lambda_0 = 3 \pm 0.5$ mm along the shock-bounded layer. The error bars were estimated by considering an average dept-of-field uncertainty ($\pm 200 \mu\text{m}$) from a pre-shot image of a known thin object. The amplitude growth model is derived by only considering conversation laws applied to the shock-bounded layers formed from colliding flows and adjusted for a single accretion shock.

The results suggest that the observed growth is compatible with the one expected from the linear model for thin-shell instability when applied to a single plasma inflow. This linear model, however, does not tell whether the layer is naturally unstable to the observed wavelength, instead only provides a possible explanation for the perturbation growth profile. In order to address this question, is required to look at the solutions of the dispersion relation for thin-shell instability.

For simplicity, we will refer to the analysis presented in (Drake, 2018) where various Vishniac-Ryu regimes of unstable shock-bounded layers are explored, providing solutions for the dispersion relation together with a more in-depth stability analysis. In the present experiments, we focus only on the standard regime of Vishniac-Ryu decelerating shock with the internal density scale length

$L\rho = -\ln(\partial\rho/\partial x)^{-1} > 0$ equal to h . Based on the already presented plasma parameters, the value of the characteristic isothermal scale height is $h = c_{s,2}^2/a \simeq 0.4$ cm and the interior density scale length, $L = h \simeq 0.4$ cm for a uniformly decelerating shock-bounded layer. Under these conditions, the stable

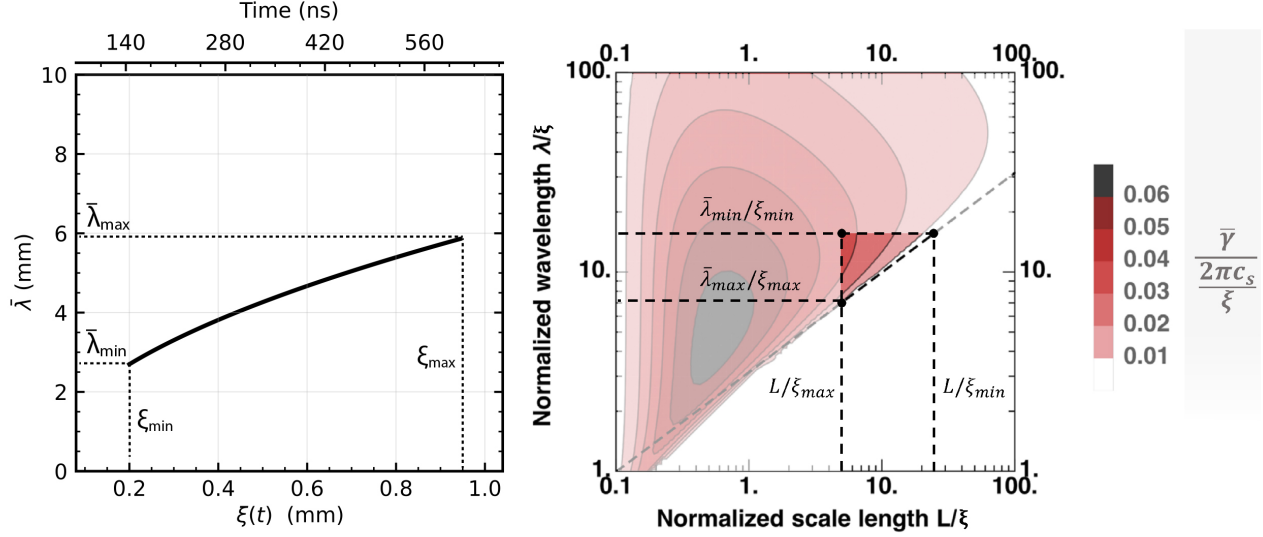


Figure A.6: (Left) Minimum unstable wavelength as a function of the thin-shell thickness, $\bar{\lambda} = \xi \pi \sqrt{h \xi}$. (Right) Contours of normalised growth rate, for $L = h$ and variations in normalised wavelength, λ/ξ and normalised (positive) scale length, L/ξ . The dashed line represents the minimum unstable wavelength. Annotated and reprinted from [37] with permission from © The American Astronomical Society.

region is given to be $\lambda < \pi \sqrt{h \xi}$ according to the solution presented in [37] for the spatial parameter $\Pi < 0$ with the minimum unstable wavelengths defined as $\bar{\lambda} = \xi \pi \sqrt{h \xi}$. Since the thickness of the layer, $\xi(t)$, changes over the duration of the experiment, a range of minimum unstable wavelengths can be computed as shown in Fig. A.6. In this regime and for the range of measured layer thicknesses, the shortest unstable modes vary between 3 mm and 6 mm which is consistent with the estimated wavelength from the experiment ($\lambda_0 = 3 \pm 0.5$ mm).

By considering the two extreme limits of possible growing modes, we can constrain the expected growth rate of perturbations for different combinations of L/ξ and λ/ξ . If we assume the expansion of the layer thickness as the primary changing parameter over time, the normalized growth rate can vary between 0.01 and 0.03 for all possible combinations of L/ξ and λ/ξ . This corresponds to a growth time ranging between $682 \text{ ns} < \tau < 2000 \text{ ns}$ for the minimum unstable wave $\lambda_{\min}(\xi_{\max})$ at the maximum thickness and $151 \text{ ns} < \tau < 454 \text{ ns}$ for the minimum unstable wave $\lambda_{\min}(\xi_{\min})$ at the minimum thickness.

These estimates provide a range for the growth time of perturbations allowed to develop under the given conditions of the Vishniac-Ryu regime and the measured thickness of the thin layer. The results of the analysis suggest that during the early phase of layer expansion, waves can be enhanced and exhibit growth in amplitude within time scales that are compatible with the duration of the experiments and with wavelengths observable by an optical self-emission camera. However, it is important to note that the system is highly dynamic, and it is possible that any of these waves could not grow rapidly enough to disrupt the layer. It is also possible that the perturbation growth slows down as less and less material reaches the reverse shock due to the finite current pulse and the limited mass of the wire array.

In conclusion, given the uncertainties associated with the assumptions made and the lack of direct measurements of the shock-bounded layer, drawing strong conclusions about the nature of this perturbation remains challenging, however, a few high-level considerations can still be made:

- Both the linear thin-shell model and the Vishniac-Ryu analysis indicate that the system can develop unstable waves based on reasonable plasma parameters for the reverse shock structure. The longest observed mode has a wavelength of $\lambda_0 \simeq 3 \pm 0.5$ mm, and its amplitude increases over time.
- The layer is not disrupted and maintains a quasi-two-dimensional profile for the entire duration of the experiment, indicating two possible scenarios: either the growth rate is too slow and only the linear phase of the instability is captured, or the driving ram pressure of the incoming plasma flow becomes less dominant over time, leading to stop any perturbation growth.
- Thin-shell instabilities can occur in accretion shock experiments, however, the complex nature of the system requires additional experimental investigations and numerical studies.

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