High Energy Density Physics
In Cluster Media

Stefan Ian Olsson Robbie
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Laser Consortium
Quantum Optics and Laser Science Group
Physics Department
Imperial College London
Prince Consort Road, London SW7 2BW
United Kingdom
Declaration of Originality

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“Success is the ability to go from one failure to another with no loss of enthusiasm.”

W.S.C.
Abstract

Gases comprised of atomic clusters have in the past been shown to exhibit extremely strong absorption of high-intensity laser pulses. By using this target medium, it is possible to use laser systems with only modest energies to create High Energy Density Plasmas. Not only are the plasmas created in this way of interest in themselves, but when properly designed, these experiments can be used as a platform for Laboratory Astrophysics studies of radiative blast waves. This thesis describes experiments which investigate the evolution of radiative blast waves, the interaction of relativistic laser pulses with large atomic clusters and the nature of the post laser-cluster interaction upstream medium into which the shock propagates.

Experiments were carried out to diagnose the properties of the upstream medium into which radiative shocks launched by the laser-cluster interaction propagate. This experiment was conducted using the Blackett Laboratory Laser Consortium Nd:Glass laser system with a novel perpendicular heating beam geometry. By introducing a time delay between the perpendicular beams, it was possible to track the propagation of a ballistic cluster disassemble wave. This wave was shown to be the product of $\sim 200$ keV ions ejected by the initial laser cluster-interaction.

Also discussed in this thesis are the results of the first laser-cluster experiment to be conducted on the Central Laser Facility’s Astra-Gemini system. Here the interaction of large atomic clusters with relativistic laser pulses is investigated. X-Ray pinhole camera images have been captured of the early time plasma created by the laser-clusters interaction. For the first time the absorption properties of large atomic clusters irradiated by a femtosecond high energy, $\sim 14$ J, laser pulse have been studied. Furthermore, the temporal evolution of radiative blast waves launched from the laser-cluster interaction is described. In the past the Vulcan laser system at RAL was used to launch blast waves which displayed velocity domain oscillations driven by the radiation emitted by the blast wave. This instability has again been observed in the work reported here and the threshold for onset has been investigated.
Acknowledgements

I should start by stating that I am a very lucky man. I have a great number of colleagues and friends that have supported me on the long journey to this point. If I have missed you in the following text, I am sorry and you should know that I am grateful to you despite the omissions made by my memory.

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¹The best office-mate in the group! At least that’s what he tells me.
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Oh and you, well you know already.

I love you all!

Stockholm, November 2013
## Commonly used Symbols

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<tr>
<th>Physical Quantity</th>
<th>Symbol</th>
<th>Value (SI)</th>
</tr>
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<tbody>
<tr>
<td>Boltzmann constant</td>
<td>$k_B$</td>
<td>$1.38 \times 10^{-23}$ J K$^{-1}$</td>
</tr>
<tr>
<td>Elementary charge</td>
<td>$e$</td>
<td>$1.60 \times 10^{-19}$ C</td>
</tr>
<tr>
<td>Electron mass</td>
<td>$m_e$</td>
<td>$9.11 \times 10^{-31}$ kg</td>
</tr>
<tr>
<td>Proton mass</td>
<td>$m_p$</td>
<td>$1.67 \times 10^{-27}$ kg</td>
</tr>
<tr>
<td>Speed of light in vacuum</td>
<td>$c$</td>
<td>$3.00 \times 10^8$ m s$^{-1}$</td>
</tr>
<tr>
<td>Permittivity of free space</td>
<td>$\epsilon_0$</td>
<td>$8.85 \times 10^{-12}$ F m$^{-1}$</td>
</tr>
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<td>Stefan-Boltzmann constant</td>
<td>$\sigma_{SB}$</td>
<td>$5.67 \times 10^{-8}$ W m$^{-2}$ K$^{-4}$</td>
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<table>
<thead>
<tr>
<th>Symbol</th>
<th>Quantity</th>
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<tbody>
<tr>
<td>$\omega_L$</td>
<td>Laser (angular) frequency</td>
</tr>
<tr>
<td>$R$</td>
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Chapter 1

Introduction.

There exist many large physical systems in the universe with complex dynamics and high energy densities. To study such phenomena experimentally, a high energy density must also be achieved in the laboratory. The advent of the modern high power laser has, over the last twenty years, given researchers access to regimes which would previously only have been found in supernova, the cores of planets or in thermonuclear explosions, which they wished to study. In these new high energy density experiments it has become possible to study, on the laboratory scale, the dynamics which control the evolution of much larger astrophysical systems. Since these laboratory experiments are repeatable, full scans of the parameters which control these systems are possible. Furthermore, the great size of many high energy density astrophysical systems is such that their evolution takes hundreds if not thousands of years, making them impossible to study observationally. Since the physical processes that govern the behaviour of the laboratory and astrophysical systems can be the same, it is possible to build experiments that, under suitable hydrodynamic conditions, are directly equivalent. This allows us to perform “scaled” experiments that go from \( \sim \) mm spatial and \( \sim \) ns temporal scales in the laboratory to thousands of kilometres and thousands of years in nature.

Typically such experiments have been carried out using only “national facility” scale laser systems with \( \sim \)100 J of energy and low shot rates [1]. However, by using the unique laser energy absorption properties of clustered gases, it is possible to drive such experiments using table top university based laser systems or indeed high repetition rate facility lasers with tens of Joules of energy [2, 3, 4, 5, 6]. This thesis is concerned with work in three linked areas: proving the viability of clustered gas media for laboratory astrophysics experiments with strongly radiative shocks, with
radiative shock experiments themselves and, for the first time, demonstrating high energy density physics experiments conducted at a high shot rate, making parameter space scans possible. This work was pursued using both smaller, university scale lasers and national facility scale systems to gather large data sets with a broad suite of diagnostics.

1.1 Motivation: Systems of Interest in Radiative Shock Laboratory Astrophysics Studies.

There exist a broad range of astrophysical systems which cannot be studied observationally, however laboratory experiments which are governed by the same physical processes can be studied. There exists a set of scaling laws, described in Sec. 2.3.4, which can be used to determine whether an astrophysical system is being scaled properly to the laboratory test system.

The thin shell structures of blast waves are susceptible to a number of instabilities and overstabilities, for example the Rayleigh-Taylor instability at an interface between different atomic species [7, 8, 9], the Vishniac overstability [10, 11] and the thermal cooling instability [5, 12, 13, 14]. Such behaviour can lead to density perturbations which can in turn, within for example supernova (SN) or in supernova remnants (SNR), lead to the nucleation points about which new astrophysical bodies can grow [15, 16, 17].

It should be noted broadly that while many experiments have been conducted by a number of groups across the world using a range of different mechanisms, there remains a fundamental problem in all of this work. Laboratory astrophysics experiments have yet to create conditions which match those set out by the scaling laws for a given system, these scaling laws are described in Sec. 2.3.4 [20]. Therefore current experiments cannot be said to be equivalent to the astrophysical systems they are intended to mimic. While this remains an issue it should also be noted that experiments are improving and moving closer to a proper match. However, it is this issue rather than that of creating particular behaviours which is the greatest challenge for the laboratory astrophysics community.

The work in this thesis is primarily intended to develop a laboratory astrophysics platform rather than answers the fundamental challenge of proper scaling at this time. A deeper discussion of the Vishniac overstability and the thermal cooling instability is provided here as an example of two phenomena which we would seek to
study with a fully developed laboratory astrophysics platform using the laser-cluster interaction to launch the shocks.

One of the primary consequences of radiation is thinning of the shock to a compression ratio greater than 4, above the limit in the non-radiative case see Sec. 2.3.2. Such thin shell shocks are more susceptible to unstable behaviour than thicker shocks [18]. Furthermore the density of the shock and its temperature strongly effect the radiative flux of the shock. If we consider the Vishniac overstability where a perturbation in the shock velocity has been introduced on an initially flat shock front we create a wave like structure of peaks and troughs, shown in Fig. 1.1 is a schematic of such a scenario [10].

![Figure 1.1: Schematic of the Vishniac overstability. (A) shows an idealised shock and (B) a shock with a perturbation on it. The mismatch in the direction of the thermal pressure behind the shock and the ram pressure on the front of the shock can be clearly seen.](image)

The pressure acting on the front of the shock is the ram pressure, due to collisions with the upstream material, which is always parallel to the direction of propagation. On the other hand behind the shock the pressure is thermal and therefore normal to the shock at each point on its surface. Therefore the pressure normal to the front of the shock falls as the shock front is no longer perpendicular to the direction the ram pressure. The thermal pressure behind the shock is, however, isotropic and so there is a net force squeezing the trough closed and then outward forming
a peak. This can result in an oscillating wave front. However as the shock velocity and compression change the radiative flux from the shock also changes, the result of this is a modulation in the preheat of the material upstream of the shock. This preheat changes the properties of the shock as it propagates through the medium and can act to amplify the instabilities present on the shock. The evolution of such a system in space is “slow” and our understanding limited by its observational nature, for example the astronomer must be observing the right star at the moment it goes supernova in order to track its’ evolution and this can only be achieved by trying to observe many stars at low resolution. Additionally even such observations can only capture the early time behaviour of the system since the evolution takes many thousands of years. Furthermore these interactions are driven by complex non-local processes and are difficult to model without high quality benchmarking data. Laboratory astrophysics studies are capable of providing such benchmarking data, and so can aid the development of models which explain the birth of new astrophysical bodies.

The thermal cooling instability (TCI) has been predicted to occur in SNR. It is a phenomenon where the changing temperature of the post shock region, cooled by the radiation losses in the system, causes the shock velocity to oscillate [12]. In a simple picture the pressure driving the shock through the ambient medium changes as the temperature of the post shock region changes. The TCI oscillation then is driven by the dependence of the cooling function on the temperature of the post shock region as:

\[ \Lambda(T) \propto T_e^\beta \]  

(1.1)

Where \( \Lambda \) is the cooling function, \( T_e \) the electron temperature and \( \beta \) the power law exponent which is a function of the temperature. When \( \beta \) fulfils this condition:

\[ \beta = \frac{d \ln \Lambda}{d \ln T} < 1 \]  

(1.2)

then an oscillation in the shock velocity should set in [12]. A single oscillation of this type in a SNR is predicted to take 500-10,000 years to occur, due to this long period TCI is an ideal candidate for laboratory radiative shock studies [19]. Furthermore modelling of the TCI is challenging since it requires accurate knowledge of the opacity of the medium and must also deal well with the changes which the radiation emitted from the blast wave drives in the upstream medium before the arrival of the shock.
The blast waves which are studied in this thesis have been shown to produce significant radiative precursors, see Sec. 6.3.1. Due to the ease of probing these blast waves they provide an ideal system to study with the aim of eventually developing well characterised radiative laboratory astrophysics experiments. In particular the TCI has already been identified in similar experiments in the past and we attempt to repeat that measurement in the work contained in this thesis [13].

1.2 Organisation of thesis.

Unless otherwise stated all equations in this thesis are written in SI notation. However, when appropriate, SI units are not used in favour of more ‘practical’ units. In these cases the units of choice will be identified by the use of Square Brackets, for example T[ns] would be period in nano seconds. Furthermore, vectors will be identified by using bold text \( \mathbf{x} \) versus scalars which will be written plainly as \( x \).

1.3 Contributors to the work and publications.

The first experiment described in this thesis was conducted on the Blackett Laboratory Laser Consortium Nd:Glass laser system, described in Chap. 4, using a Parker 99 valve. The set-up and results of this experiment are discussed in Chap. 5. This experiment and analysis of the data was conducted by H.W. Doyle and the author.

The second experiment described in this thesis was performed at the Central Laser Facility (CLF) on the Astra-Gemini laser system, where atomic clusters were irradiated at relativistic intensities. This experiment was supported by the CLF Target Fabrication group, the Astra-Gemini laser operators and target area link scientists. The experimental team was drawn from Imperial College, the University of Oxford, the CLF and the Atomic Weapons Establishment. The team comprised, in no particular order, of D. Bigourd, S. Patankar, H.F. Lowe, K. Mesceki, C. Price, H.W. Doyle, J. Fryth, E.T. Gumbrell, D.R. Symes, R.A. Smith and the author. At the time of writing, modelling of the launched shocks was being conducted by R. Scott of the CLF. Analysis of the data produced on this experiment was conducted by H.F. Lowe, D. Bogourd and the author. The data pertaining to the shocks launched during this experiment are described in Chap. 6 and data pertaining to the early time plasma and laser-cluster interaction is described in Chap. 7.
The following publications and presentations are based on the work described in this thesis:


The following publication is not described in this thesis but the author did participate in the work:

Chapter 2

Theoretical Background

The purpose of this chapter is to provide a basic introduction to both the theoretical underpinnings of the experimental techniques used in this work and an understanding of the data generated by the experiments described in this thesis. This begins with a brief introduction to the definition of a plasma, how to go about the modelling of a plasma and then moves on to laser-matter interactions and shock physics.
2.1 Plasma physics concepts - an introduction.

In its simplest form a plasma is an ensemble of electrons, ions and neutral atoms or molecules residing in both internal and external electric and magnetic fields [4, 21, 22]. The electromagnetic fields are an extremely important ingredient in the behaviour of the plasma governing both local and collective behaviour and are often modified or generated by the plasma itself. Plasma is a very common form of matter in the universe at large and as such its study covers a very broad field. Since this thesis is concerned with laser produced plasmas in cluster media this section seeks to describe the most important properties to the launching of shocks in these plasmas.

2.1.1 Definition of a plasma and the Debye length.

We consider a “simple” idealised plasma containing $n_e$ electrons and $n_i$ ions per unit volume with the ions in a variety of charge states denoted as $Z_i^*$. The electrons and ions will tend to group together, due to their opposite sign charges, reaching a minimum separation determined by their thermal energy. From this arises the concept of the Debye length, $\lambda_D$, the characteristic length scale of the plasma system. The Debye length is calculated as shown here:

$$\lambda_D^{-1} = \sqrt{\frac{e^2}{\epsilon_0 k_B T_e} \left[ \frac{n_e}{T_e} + \sum_i \frac{Z_i^*}{T_i} \right]}$$ (2.1)

Where $e$ is the electron charge, $\epsilon_0$ is the permittivity of free space, $k_B$ is the Boltzmann constant, $T_e$ & $T_i$ are the electron and ion temperatures respectively, $n_e$ & $n_i$ are the electron and ion number densities respectively and $Z_i^*$ is the ion charge state, over which the sum runs. Electron motion generally occurs on a time scale so short when compared to the ions, that the ions can be considered “frozen” into the background. Using this fact we can approximate the Debye length for the electrons to $\lambda_D = \sqrt{\frac{e^2 n_e T_e}{k_B n_e^e}}$, in this thesis the references to the Debye length will refer to the electron Debye length. If there is a charge balance between the electrons and ions ($n_e = \sum_i Z_i n_i$) over the Debye length, then the plasma is referred to as quasi-neutral.

A better way of understanding the physical nature of the Debye length is to consider a test charge, $q_{test}$, in a plasma where the total charge density can be expressed as:
\[ \rho_c = \rho_{\text{test}} - e n_e \exp\left(\frac{e \phi}{k_B T_e}\right) + \sum_i e n_i Z_i^* \exp\left(\frac{-e Z_i^* \phi}{k_B T_i}\right) \]  

(2.2)

Where \( \rho_c \) is the total charge density, \( e \) is the electron charge, \( n_e \) and \( n_i \) are the electron and ion number densities, \( T_e \) and \( T_i \) are the electron and ion temperatures, \( Z_i^* \) is the ion charge state, \( k_B \) is the Boltzmann constant and \( \phi = \phi(x, y, z) \) the scalar electric potential in the plasma.

In the so called weak coupling limit, where the thermal energy of the plasma is greater than the electric potential, \( 1 \gg \frac{e \phi}{k_B T_e} \), Equation 2.2 becomes \( \rho_c = \rho_{\text{test}} + \phi \epsilon_0 \lambda_D^2 \).

This takes the form of Poisson’s equation, \( \nabla^2 \phi = p_c / \epsilon_0 \), which when solved leads to:

\[ \phi(r) = \frac{q_{\text{test}}}{4\pi \epsilon_0 r} \exp[-r/\lambda_D] \]  

(2.3)

As is clear from Eq. 2.3 the potential of a single charge within a plasma falls away very quickly. When compared with the slower fall in potential of a charge in vacuum, \( \phi \propto 1/r \), it is clear that this is one of the defining characteristics of a plasma. This is known as Debye Shielding and for this to hold, the plasma must satisfy the Plasma Parameter conditions:

\[ n_e^{-1/3} \ll \lambda \ll L \]  

(2.4)

Where the plasma parameter is the number of electrons contained in a Debye Sphere, a sphere with radius \( \lambda_D \). For a “genuine” plasma the plasma parameter must be sufficiently large, such that \( N_D = \frac{4\pi}{3} n_e \lambda_D^3 \gg 1 \). This can alternatively be stated as \( \lambda_D \gg n_e^{-1/3} \) and in an ideal plasma \( N_D \) goes to infinity.

We can use the argument in Eq. 2.4 to define a simple set of conditions which define an ideal plasma. Which will be in the weak coupling limit and quasi-neutral as shown in Eq. 2.4, where \( L = n_e / |\nabla n_e| \) the so called plasma scale length. If we imagine a plasma typical of a laser-cluster interaction in hydrogen gas with \( T = 1 \) keV and \( n_e = 10^{19} \text{ cm}^{-3} \) it would lead to \( n_e^{-3} \approx 5 \text{ nm} \), \( \lambda_D \approx 50 \text{ nm} \) and \( L \sim 100 \mu\text{m} \). This is reasonably close to an ideal plasma but it should be noted that many systems that are only weakly ionised are sometimes referred to as plasmas even though they do not fully match the condition set out in Eq. 2.4.
2.1.2 Description Of Plasmas.

The simplest description of the motion of a plasma is to consider all the constituent particles and their motion through an external electromagnetic field. The equation governing a given particle is simply the Lorentz force on that particle. This is given by:

\[
m_j \frac{d\mathbf{v}_j}{dt} = q_j \left( \mathbf{E} + \mathbf{v}_j \times \mathbf{B} \right) + \mathbf{F}
\]

(2.5)

Where \( q, m \) and \( \mathbf{v} \) are the charge, mass and velocity of the particle in the field with \( \mathbf{E} \) the electric field, \( \mathbf{B} \) the magnetic field and \( \mathbf{F} \) any additional forces for a given particle \( j \). While this description is entirely accurate, it does require huge computational power to apply in any case with more than a handful of particles. Therefore, when actually modelling “real” plasmas with many constituents, a more sophisticated approach is required.

A significant simplification of this model can be accessed by applying a statistical approach using the Boltzmann equation. This is a kinetic description where it is assumed that all collisional contributions to the electromagnetic field are shielded within the Debye sphere, leading to a collisionless Boltzmann equation. This is called the Vlasov Equation as shown here:

\[
\frac{\partial f}{\partial t} + \mathbf{u} \cdot \nabla f + \frac{q}{m} (\mathbf{E} + \mathbf{u} \times \mathbf{B}) \cdot \frac{\partial f}{\partial \mathbf{u}} = 0
\]

(2.6)

Where \( f(\mathbf{r}, \mathbf{v}, t) \) is the phase space distribution function of the plasma, \( \mathbf{u} \) the fluid velocity and \( m \) the particle mass.

Eq. 2.6 is written for every species in the plasma which you wish to model. However, this description is still concerned with the relatively fine scale behaviours of the system, the work described in this thesis is more concerned with the macroscopic behaviour of certain plasmas and as such a fluid description of the plasma is often more appropriate. By taking velocity moments of the Vlasov equation and assuming an isotropic pressure distribution, we arrive at the Euler Fluid Equations [22]:

\[
\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{u}) = 0
\]

(2.7)

\[
\rho \left( \frac{\partial \mathbf{u}}{\partial t} + (\mathbf{u} \cdot \nabla) \mathbf{u} \right) = \rho c (\mathbf{E} + \mathbf{u} \times \mathbf{B}) - \nabla p
\]

(2.8)
\( \frac{\partial p}{\partial t} + \nabla p = -\gamma p \nabla \cdot \mathbf{u} \)  \hspace{1cm} (2.9)

Where \( \rho, \rho_c \) and \( p \) are the density, charge density and pressure respectively. The “continuity equation”, Eq. 2.7, describes the conservation of mass in the system. The ‘equation of motion’, Eq. 2.8, describes the change in momentum density as a result of the electromagnetic forces and the pressure gradients in the systems. The “energy equation”, Eq. 2.9, describes the conservation of energy in the system. The system is assumed to be polytropic, where the pressure is related to the density according to \( p \propto \rho^\gamma \). Here \( \gamma \) is the Polytropic Index of the system and, if given by the ratio of specific heat at constant pressure to specific heat at constant volume, \( \gamma = c_p/c_v \) it can be referred to as the Adiabatic Index. Further, according to thermodynamics for an ionised classical ideal gas, with \( T_i \ll T_e \), \( \gamma \) can be expressed based on the number of degrees of freedom in the system such that, typically, \( \gamma = \frac{\nu+2}{\nu} = 5/3 \).

As the temperature of the plasma increases and the density drops, the collisional coupling between the ions and electrons decreases, eventually reaching a state where the ions and electrons can be treated as separate fluids. This leads to the Two-Fluid model of the plasma where two separate phase space distributions are used for the electrons and ions respectively. Additionally, it is often the case that plasmas will contain a small number of ions and/or electrons that are much more energetic than the primary populations. In such cases the model can be extended by including some extra ‘fluids’ to represent these secondary populations with their own phase space distributions.

One additional component is required to complete the two-fluid description: Maxwell’s Equations. Maxwell’s equations relate the electric and magnetic fields to the charge and current density of the plasma, as shown here:

\[
\nabla \cdot \mathbf{E} = \frac{\rho_c}{\epsilon_0}, \quad \nabla \cdot \mathbf{B} = 0
\]
\[
\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t}, \quad \nabla \times \mathbf{B} = \mu_0 \mathbf{J} + \frac{1}{c^2} \frac{\partial \mathbf{E}}{\partial t}
\]  \hspace{1cm} (2.10)

Where \( \rho_c \) is the charge density, \( \mathbf{J} \) is the current density, \( \mu_0 \) is the permeability of free space and \( \epsilon_0 \). This addition completes the two-fluid model but has assumed that the system is collisionless. If we wish to extend the model to a collisional system then 2.6 must be set equal to \( (\frac{\partial f}{\partial t})_{\text{coll}} \) rather than to zero. Where \( (\frac{\partial f}{\partial t})_{\text{coll}} \) is the rate
of change of the phase space distribution for the electrons and ions. A remaining
issue is that by taking moments, each subsequent equation becomes reliant on its
predecessor, with the 1st depending on the 0th and 2nd on 1st etc, leading to an
infinite set of equations. We can truncate this process by making certain limiting
assumptions about heatflow, leading to the Equation of State (EOS). If we assume
rapid heat flow in the system then we achieve a constant temperature throughout
the system. This “closes” the loop leaving a complete description of the plasma [22].

2.2 Laser-Matter Interactions.

When focusing a high intensity laser onto matter the interactions are typically com-
plex and dependent on many variables. In this section the author will explore the
primary factors to consider in a laser-matter interaction, and particularly the in-
teraction of a laser with an atomic cluster, the target medium of choice in the
experiments described later in this thesis.

2.2.1 Targets In Laser Matter Interactions.

A broad variety of target materials and formats are used in experiments which range
from low density gases (n~10^{17} \text{cm}^{-3}) through aerogels (n~10^{19} \text{cm}^{-3}) to bulk solids
(n~10^{23} \text{cm}^{-3}). Not only do the different densities change the physics that govern
the plasma but they also alter the coupling of laser energy into the target. Fig.
2.1 shows a variety of commonly used targets covering the range of scales used in
laser-matter interactions.

Bulk solid target interactions with laser fields have been studied since the advent
of the laser in the 1960s. They have been used to study systems ranging from X-ray
generation [23], fusion research [24] to laboratory astrophysics and EOS measure-
ments [25]. Due to the thickness of such targets, the laser will not directly heat the
material beyond the skin depth and the energy coupling at the surface can be rea-
sonably efficient, but heating deep into the material is less efficient. For this reason
experiments that require bulk heating, such as EOS studies, rely often on different
indirect heating sources, which are themselves often laser driven, for example proton
heating [6, 26, 27] or x-ray heating [28].

Thin foils are a popular target for the generation of protons or ions which can be
accelerated by the Target Normal Sheath Acceleration (TNSA) mechanism. Here a
high intensity laser is focused onto a foil, generating a plasma and space charge on
the back surface of the foil. This space charge drags ions along with it accelerating them out of the back of the foil toward a secondary target. For further details see [27, 29].

Microdroplet targets are typically of a size greater than the skin depth of the laser. As a result of this, the surface of the microdroplet is heated directly by the laser and then, due to the still relatively small size of the microdroplet, the heat is rapidly transported through the bulk material. However, since the microdroplet is still of the order of the wavelength of the laser, electric field enhancements can occur on the surface, complicating the dynamics of the system. This enhancement has been shown to increase the energy coupling of the laser into the microdroplets, leading to enhanced x-ray production [30].

Atomic clusters consist of aggregates of atoms from a few tens of atoms up to ~10,000 atoms held together temporarily by weak inter atomic, molecular or metallic binding forces, depending on the species. Laser energy in short pulses (<10 ps) is very well absorbed(<90%) by clusters, while also allowing good penetration into the medium. The experiments described in this thesis make extensive use of cluster laser absorption properties and will be described in more detail in Section 2.2.4.

Molecular and atomic gases do not absorb laser energy particularly well (<5%), however they do make it possible to drive a broad range of experiments. For example using High Harmonic Generation (HHG) techniques it is possible to map out the orbitals of atoms and indeed to investigate the bonds that make up molecules [31]. Further, HHG can also produce coherent beams of such broad bandwidth that a
single sub-femto second pulse can be produced [32]. Both of these applications make use of ultra-short (\(\sim 10\) fs) laser pulses and are typically considered to be quantum optics studies rather than plasma physics experiments, although plasma physics may play a key role, in particular limiting the phase matching due to the dispersion properties of plasmas.

A further use of atomic gas targets is to build laser particle accelerators using relativistic laser pulses [33]. Here a plasma wave is created which traps electrons and accelerates them to extremely high energies. GeV energies have been achieved using Wakefield Accelerators which, it is hoped, will be able to replace kilometre-long conventional linear accelerators with centimetre-long accelerators [34].

### 2.2.2 Motion of Electrons In a Laser Field.

Clearly free electrons in a laser field experience an electromagnetic force due to that field. The two halves of a laser field propagating in the x direction can be described as \(E(r, t) = E_0 \cos(\omega_L t - kx)\hat{y}\) and \(B(r, t) = B_0 \cos(\omega_L t - kx)\hat{z}\), where \(\hat{y}\) and \(\hat{z}\) are the unity vectors in y and z respectively, \(k = 2\pi/\lambda\) is the wavenumber and \(\omega_L = 2\pi c/\lambda\) is the laser angular frequency with laser wavelength \(\lambda\). If we combine this with the general Lorentz force, Eq. 2.5, with no additional external forces and an electron travelling at a velocity \(v_e\) with respect to the laboratory frame we obtain:

\[
m_e \frac{dv_e}{dt} = -e \left( E_0 \cos(\omega_L t - kx)\hat{y} + v_e \times B_0 \cos(\omega_L t - kx)\hat{z} \right) \tag{2.11}
\]

If we apply the third Maxwell equation, shown in Eq. 2.10, then we find that \(|E| = |B|\omega_L/k = cB_0\) with \(\omega_L/k = v_{phase} = c\) where \(v_{phase}\) is the phase velocity of the light. From this it is clear that \(E \gg vB\) if \(|v_e| \ll c\) and hence we need only consider the contributions of the electric component of the laser field for sub-relativistic particle energies. If we rearrange Eq. 2.11 and integrate for time we arrive at:

\[
v_e = -\frac{e^2 E_0^2}{m_e \omega_L} \sin(\omega_L t - kx) \tag{2.12}
\]

From this equation we can see that the electron oscillates in the laser field. From this oscillation we arrive at the Quiver Velocity which is given by the maximum velocity of the electron oscillation \(v_{quiver} = \frac{eE_0}{M_{\text{eql}}}\). We can work out the quiver energy by substituting into the kinetic energy equation and averaging over a laser full cycle which leads to:
\[ \xi_{\text{quiver}} = U_p = \frac{1}{2} m_e \langle v_e^2 \rangle = \frac{e^2 E_0^2}{4 m_e \omega_L^2} \sim I_L \lambda^2 \]  

(2.13)

Where \( I_L \) is the laser intensity and \( \lambda \) is the laser wavelength, this is more commonly known as the Pondermotive Potential \( (U_p) \). It is possible to estimate the Pondermotive potential of the laser field as \( U_p[\text{KeV}] \approx I_L[10^{18}\text{Wcm}^{-2}] (\lambda[\mu\text{m}])^2 \). Also we can use \( U_p \) to obtain the Pondermotive Force, via \( \mathbf{F}_p = -\nabla U_p = -\frac{e^2}{4 m_e \omega_L^2} \nabla E_0^2 \) which describes the force on an electron from a laser field. This force will act on the electrons in the laser field to push them away from the higher intensity regions normal to the propagation direction of the laser pulse. This of course means that in, for example, a Gaussian beam profile, an electron will eventually exit the laser field and propagate away having acquired a velocity \( v \sim v_{\text{quiver}} \).

Whilst this description is valid when electrons are travelling at non-relativistic speeds, it breaks down in the experiments described in this thesis. The laser intensities used for some of the experiments described in this thesis are \( >10^{18}\text{Wcm}^{-2} \) a point where the velocity of the electron oscillating in the laser field approaches the speed of light, \( c \). When this is the case the magnetic field term in the Lorentz force, \( \mathbf{v}_e \times \mathbf{B} \), is no longer negligible as compared to the electric field contribution. Additionally, at near relativistic speeds, the electron mass will be altered from its rest value, modifying Equations 2.12 and 2.13.

The magnetic field contribution results in an additional oscillating drift motion. If we observe the movement of the electron from a frame moving at the electron drift velocity, a figure of eight motion will be seen. The energy of this system must now be described relativistically based on the electron velocity, \( v_e \), as \( E_e = \gamma m_e c^2 \). Here \( \gamma \) is the Lorentz factor, which as usual is given by \( \gamma = 1/\sqrt{1 - v_e^2/c^2} \). If we modify the quiver velocity for the relativistic case we obtain \( v'_{\text{quiver}} = \frac{e E_0}{\gamma m_e \omega} \) which we can use to arrive at the Lorentz factor for this case:

\[ \gamma = \sqrt{1 + \frac{e^2}{2 \pi^2 \epsilon_0 m_e^2 c^5} I L \lambda^2} = \sqrt{1 + a_0 I_0} \]  

(2.14)

Where \( I_0 \) is the laser intensity that corresponds to maximum quiver velocity and \( a_0 \) is the normalised vector potential. The normalised vector potential is given by \( a_0 = eE/m_e c \omega_0 \) where \( \omega_0 \) is the laser angular frequency, if \( a_0 = 1 \) then the electron has the same kinetic energy as its rest mass and when \( a_0 > 1 \) the interaction can be considered relativistic [35].

It is also possible to write Eq. 2.14 in more practical units such that \( \gamma \approx \)
(1 + \lambda / 1.4 \times 10^{18} \text{Wcm}^{-2})^{1/2}$. From this it is clear that when the intensity of the laser exceeds $1.4 \times 10^{18} \text{Wcm}^{-2}$ then the system becomes relativistic. This coupling between the laser field and plasma is dependant on the laser field penetrating into the plasma. When the plasma reaches the critical density, the laser pulse will be reflected from the surface. A plasma that is below this threshold, such that the laser pulse can penetrate it, is said to be underdense, whereas a plasma above this threshold is said to be overdense.

2.2.3 Ionisation Mechanisms.

At low intensities the Photoelectric Effect is the dominant ionisation mechanism. Here a single photon of sufficient energy liberates an electron from its parent atom. The atom absorbs a photon with energy $E = h \nu$, where $h$ is Planck’s constant and $\nu$ is the frequency of the photon. If $E$ is large enough to overcome the binding potential of the atom then the electron is liberated and escapes with a kinetic energy of

$$E_{Km} = h(\nu - \nu_0)$$

where $\nu_0$ is the minimum photon frequency required to liberate the electron. This picture, however, breaks down when the intensity of the laser increases.

When the laser intensity exceeds $10^{10} \text{Wcm}^{-2}$ Multiphoton Ionisation (MPI) will start to occur. In this case so many photons are arriving in such a short time that an electron can be liberated from an atom by ionising from one virtual state to another until escape. Each incoming photon will elevate the electron energy by $h \nu$. However, the liberated electron will receive very little kinetic energy in this picture. At even higher intensities, $\geq 10^{14} \text{Wcm}^{-2}$, enough photons are incident on the atom that not only is the electron liberated from the atom, but also significant kinetic energy is imparted. Here the electron will escape fully into the continuum. This process is known as Above Threshold Ionisation (ATI).

In very short intense laser fields the atomic potential is no longer simply perturbed but is instead significantly distorted. The laser field can skew the potential of the atom, increasing the probability of Tunnel Ionisation (TI) occurring or at even more extreme levels the potential can be suppressed below the energy level of the electron, so called Barrier Suppression (BS). The intensity at which barrier suppression occurs is $I_{BS} = \frac{\pi^2 \alpha^2 V_I^4}{2Z^2 e^6 \epsilon}$ [4], where $V_I$ is the electric potential of the atom.

It is possible to determine whether the MPI or the BS regime is dominant for a
given atomic species and laser intensity using the Keldysh Parameter \((\Gamma_k)\):

\[
\Gamma_k = \omega_L \sqrt{\frac{2m_e V_I}{e E_L}} = \sqrt{\frac{V_I}{2U_p}}.
\]  

(2.15)

If \(\Gamma_k \gg 1\) then MPI is the dominant ionisation mechanism in the field and clearly if \(\Gamma_k \ll 1\) then TI will be dominant. However, it should be noted that the Keldysh parameter is too strict a rule, and both ionisation processes will be present over a broad range of parameters. This is particularly true when considering a whole laser pulse where the intensity will rise and then fall after the peak giving different intensities at different times, together with spatial variations in intensity across the laser focus.

2.2.4 Laser-Cluster Interactions.

Clusters become very strong absorbers, >90%, when an intensity threshold is broken, typically this intensity will be around \(\sim 10^{12}\) Wcm\(^{-2}\) [4, 36, 37, 38, 39, 40]. However, clusters are very complex and no theory or model successfully describes their complete behaviour over the full range of laser or cluster parameters. Ditmire et. al. developed the Nanoplasma model which, over a limited range of experimentally useful conditions, is able to predict much of the observed cluster behaviour described in this thesis [41]. However the nanoplasma model cannot be applied to the results in Sections 6 and 7 since the laser pulses involved are too short, and the clusters too large for its simplifying assumptions to be valid.

In the nanoplasma model the clusters are treated as spheres with uniform density and temperature. The cluster must also have a radius, \(R_C\), which holds with \(R_C \gg \lambda_D\). Furthermore, to maintain the uniform temperature across the cluster \(R_C\) must be smaller than the skin depth which is the depth to which the laser field penetrates into the cluster. Energy is assumed to be coupled from the laser field into the cluster via collisional inverse Bremsstrahlung as if the cluster were an over dense plasma.

If \(R_C < \lambda_{Laser}\) the sphere can treated as if it were a dielectric under a constant electric field given by \(E_0\) then we arrive at as an expression of the electric field:

\[
E = \frac{3E_0}{|\epsilon + 2|}
\]

(2.16)

Where \(\epsilon = 1 - \frac{\omega_p^2}{\omega_L (\omega_L + \nu_{ei})}\) is the complex dielectric constant, \(\nu_{ei}\) is the electron-ion collision frequency derived from the Drude model, \(\omega_p\) & \(\omega_L\) are the plasma & laser
frequency respectively [42].

Clearly as $|\varepsilon + 2| \rightarrow 0$ a maximum in the electric field will occur in Eq. 2.16 when $\omega_L = \omega_p/\sqrt{3}$, the so called Mie Frequency of the cluster, this is the point at which the cluster is in a resonance. Alternatively, and more usefully, this condition can be expressed in terms of the internal electron density of the clusters and the critical electron density of the laser field which yields $n_e/n_{\text{Critical}} = 3$. However, it should also be noted that when $n_e/n_{\text{Critical}} > 6$ a shielding effect starts to occur within the cluster.

The cluster undergoes expansion driven either by hydrodynamic pressure from the hot electrons within the cluster or Coulomb pressure resulting from positive charge build up in the cluster caused by escaping fast electrons. The dominant mechanism can be determined by comparing the temporal duration of the laser field to the disassembly time of the cluster, given by $t_{\text{Dis}} \approx 0.8 \sqrt{\frac{4\pi\epsilon_0 m_i n_i e^2}{n_e}}$; if the field is similar to or shorter than $t_{\text{Dis}}$ then the Coulomb pressure will dominate since a large portion of the electrons will escape before significant expansion has occurred. However this does also rely on the pondermotive potential, $U_p$, of the laser field being sufficiently large to liberate the electrons [43]. In the hydrodynamic case a number of electrons have been optical ionised early in the laser field out of their parent ions, these electrons are then accelerated by the main laser field which transfers energy to the cluster via inverse Bremsstrahlung [41].

The hydrodynamic pressure is given by $P_{\text{Hydro}} = n_e k_B T_e$ and scales with the radius of the cluster as $1/R_C^3$. The Coulomb pressure can be modelled using Maxwellian free streaming electrons out of the cluster which then determines the charge build up in the cluster treated as a spherical conductor. The charge across the surface of the cluster, $Q_e$, allows the energy density per unit area of surface to be derived. This is then used to derive the Coulomb pressure as $P_{\text{Co}} = \frac{1}{2} \frac{1}{(4\pi\epsilon_0)^2} \frac{Q^2 e^2}{R_C}$ which scales with the radius of the cluster as $1/R_C^4$. Due to the strong scaling with radius the hydrodynamic pressure will dominate the majority of the evolution of the cluster rather than the Coulomb pressure.

If the laser pulse is of sufficiently long duration, then as the cluster expands the electron density will fall and the cluster will again fulfil the condition that $n_e/n_{\text{Critical}} = 3$ entering another resonance. This description is intended only as a brief introduction and further detail of the behaviour of clusters can be found in [5].
2.2.5 Overview of Previous Work Involving Laser-Cluster Interactions.

High intensity laser-cluster interactions have been studied for some time, this work was begun by the observation of anomalous x-ray line emission from xenon and krypton in high pressure gas jets by McPherson et. al. in 1994, explained using their “hollow atom” model [44]. This model was later shown to be an error by Ditmire et al, and it was instead shown that it was atomic clusters in the gas plume which resulted in the anomalous x-ray emission [41, 45]. This work was largely motivated by the desire to create useful sources of x-rays without generating debris which could damage sensitive equipment. In parallel with this work interest grew in the high energy electrons and ions which could be generated by the laser-cluster interaction. Earlier experiments focused on the generation of fast electrons of 3 keV peak energy [46, 47] and the generation of noble gas ions of 1 MeV energy [40, 48] with only modest laser intensities of $\sim 10^{16}$ Wcm$^{-3}$. This work has now evolved to the point where the clusters are being considered as an injection source for a laser wakefield accelerator [49].

Clusters have also be used as a fuel medium in laser fusion experiments [43, 50, 51]. While the clusters act as an intense neutron source it does not appear possible to scale this work to energy generation applications. These experiments have sought to use the strong energy absorption properties of the clusters along with the possibility of achieving field enhancements on smaller clusters [52]. In the late 1990s it was recognised that the plasma filament formed by the laser-cluster interaction evolved in a shock structure which could potentially be used to conduct laboratory astrophysics studies with more accessible table-top laser systems [53]. The key measurements in this field have been high spatial resolution imaging of the shock and radiative precursor [2, 54] and full time resolved trajectory measurements [55]. The high spatial resolution images of the full structure provide benchmarking points for the radiation hydrodynamics codes which are used to model radiative shocks and have demonstrated the complex effects of radiation in these systems [4]. The trajectory measurements have identified the astrophysically relevant Chevalier and Imamura temporal instability which had never before been observed in other laboratory astrophysics studies [12, 13]. Further attempts have been made to initiate the Vishniac instability with, at this time, no clear success [56]. However, they have demonstrated the possibility of using creative experimental arrangements to investigate such phenomena, arrangements which cannot easily be achieved with
2.3 Shocks, Blast Waves and Scaling Laws.

2.3.1 Shock Waves, Rarefactions and Blast Waves.

A shock wave will occur when matter is forced to move at a speed faster than that at which the ambient medium is capable of moving out of the way [18]. The simplest picture of such a system is a piston being driven, impulsively, into a closed cell from one end. The cell is filled with gas, initially at a pressure $\rho_1$, when the piston is driven inward it moves the gas immediately in front of it faster than the material further upstream can move out of the way. This results in a build-up of gas in front of the piston which, when the piston itself stops, continues to propagate further through the cell with a rarefaction wave forming behind it. This is illustrated conceptually in Fig. 2.2.

![Diagram of shock formation in a closed cell with an impulsively driven piston.](image_url)

Figure 2.2: Conceptual illustration of a shock formation in a closed cell with an impulsively driven piston (A). As the piston is driven further into the cell, gas is pushed forward into the upstream medium at such a speed that the upstream medium itself cannot move out of the way (B). This leads to a build up of gas, i.e. a density increase, in front of the piston. When the piston stops moving the high density region will continue to propagate through the cell, with a rarefaction wave forming behind it (C).
If we assume the piston to have instantaneously reached the speed $u$ and to have travelled for a time $\Delta t$ then the piston will have moved $u\Delta t$. This will result in a compression of gas in front of the piston with mass per unit area $\rho_1 u \Delta t$ which, results in a density of $\rho_2$ in front of the piston, where $\rho_2 > \rho_1$.

For a medium with a sound speed $c_s$ the material perturbation will be subsonic if $u < c_s$ and a shock would propagate with a speed $u_{\text{shock}} = c_s$. In a polytropic gas the sound speed has a temperature dependence described by $c_s = \sqrt{\gamma RT}$, where $\gamma$ is the polytropic index, $R$ is the ideal gas constant and $T$ is the temperature. If the temperature difference across this density wave is sufficiently large, with the sound speed ahead of the density wave much higher than that behind the density wave, then mass will be “swept up” to a compression front. The density profile will steepen into a step-like discontinuity, it is this discontinuity which is the shock. For a medium where $u > c_s$ the mass in front of the piston will travel so fast that a shock will necessarily have to form, in such a case the piston will have been supersonic. The Mach number of a shock is given by the ratio of the shock velocity and the sound speed of the medium it is propagating in, $M = u_{\text{shock}} / c_s$.

The shock front will propagate a distance $u_{\text{shock}} \Delta t$ during the compression time, $\Delta t$ and the total volume of compressed gas per unit area will be $(u - u_{\text{shock}}) \Delta t$. Further, mass conservation results in a density increase of $ho_2 (u_{\text{shock}} - u) = \rho_1 u_{\text{shock}}$.

If the piston is suddenly withdrawn then the density will drop, launching a rarefaction wave into the cell. This is a product of the expansion of the gas back toward the piston and if the piston continues to be withdrawn then the rarefaction wave will continue to propagate into the medium in the same fashion as the compression shock described above would. However, in the mixed case where the piston is first driven forward and then abruptly stops, a compression wave will be launched followed by a rarefaction wave. If the rarefaction wave catches up with the shock front it may form into a blast wave structure [25].

The example of an impulsively driven piston moving into a gas demonstrates well the basic physics behind several launch scenarios, for example high explosives, laser matter interactions and astrophysical shocks. In terms of modelling such a system we can conveniently decouple the energy deposition mechanism from the shock evolution. The system can be treated as two separate regions: the ambient medium and the extremely hot material generated from laser heating or an explosion. When the simulations are run, region two will expand into region one creating the shock. Given the near instantaneous heating resulting from the laser-cluster interaction on a hydrodynamic time scale, this description will inform the shock model described...
in this thesis [2, 3, 4, 5, 6].

2.3.2 Jump Conditions.

The jump conditions of a shock arise from mass, momentum and energy conservation across the shock. Consider a system with the density of some general property $Q$ given by $\rho_Q$ which is changed with a flux given by $\Gamma_Q$ such that $\frac{\partial}{\partial t} \rho_Q + \nabla \cdot \Gamma_Q = 0$. If we integrate across two regions $(X_1 & X_2)$, i.e. across the shock front, we find that as $X_1 \to X_2$ we arrive at a simple relation $\Gamma_Q(X_1) = \Gamma_Q(X_2)$. This is true for the mass, momentum and energy conservation and accordingly leads to three equations for a simple scalar case without electric or magnetic fields [18, 25]. Where the conservation in mass is given by:

$$\rho_1 u_1 = \rho_2 u_2$$ (2.17)

Where the conservation in momentum is given by:

$$\rho_1 u_1^2 + p_1 = \rho_2 u_2^2 + p_2$$ (2.18)

Where the conservation in energy is given by:

$$\rho_1 u_1 (\epsilon_1 + \frac{u_1^2}{2}) + p_1 u_1 = \rho_2 u_2 (\epsilon_2 + \frac{u_2^2}{2}) + p_2 u_2$$ (2.19)

These equations are known as the jump conditions or as the Rankine-Hugoniot relations, where $u$ is speed, $\rho$ the mass density and $\epsilon$ the internal energy per unit mass. A “Strong Shock” occurs when the Mach number of the shock is so large that only terms that contain the largest power of the Mach number need to be included for an accurate description [25]. If we consider a strong shock where we can neglect the un-shocked fluid velocity and pressure in a polytropic gas, with $\rho \epsilon = p/(\gamma - 1)$, then in the laboratory frame we arrive at these equations 2.20 [18, 25, 58]:

$$\rho_2 = \rho_1 \frac{(\gamma + 1)}{(\gamma - 1)} \quad p_2 = \rho_1 u_s^2 \frac{2}{(\gamma + 1)} \quad u_2 = u_s \frac{2}{(\gamma + 1)}$$ (2.20)

These are the jump conditions for a strong shock in an ideal gas. As is clear from these equations with an ideal gas, $\gamma = 5/3$, we arrive at a maximum compression ratio of $\rho_2/\rho_1 = 4$. This means that even for a very strong shock there is a limit as to how high a degree of compression can be achieved without radiation [18].
2.3.3 The Sedov-Taylor Solution of a Blast Wave.

In a system governed by hydrodynamics it is possible to arrive at a similarity solution, where the solution appears the same at all scale lengths for a system that has a distinguishable origin. In particular the Sedov-Taylor Solution analytically describes the “blast wave phase” of the evolution of a shock and can be applied to both the cylindrical and spherical shocks. This is the adiabatic phase when the evolution of the blast wave is only governed by the abruptly deposited energy and the ambient density of the surrounding medium. The mass swept up by the blast wave as it propagates through the medium is contained in a thin shell immediately behind the shock. The post shock material will have an extremely low mass, but will have high temperature and pressure. Using this self similarity we arrive at a solution to the time dependence of the radius given by [25, 59]:

\[ R(t) = At^\alpha \]  

(2.21)

In Eq. 2.21 \( R(t) \) is the blast wave radius, \( t \) is time with \( t = 0 \) defined as the point of energy input, \( \alpha \) is the so called deceleration parameter and \( A \) is a factor based on the initial conditions and geometry of the system. The deceleration parameter is dependent on the geometry of the system, in a spherical geometry \( \alpha = 2/5 \) and in cylindrical geometry \( \alpha = 1/2 \) in a simple system with no energy loss or gain. The experiments investigated in this thesis have a cylindrical symmetry with the shock launched from a plasma filament created by focusing a single laser pulse into atomic clusters and as such \( \alpha = 1/2 \) will be the standard solution unless otherwise stated.

In cylindrical coordinates Eq. 2.21 can be written exactly for a Sedov-Taylor blast wave as shown here [53, 60]:

\[ R(t) = \left[ \frac{4(\gamma + 1)(\gamma - 1)^2}{\pi(3\gamma - 1)} \right]^{1/4} \left( \frac{E_i}{\rho_0} \right)^{1/4} t^{1/2} \]  

(2.22)

In Eq. 2.22 \( E_i \) is the initial deposited energy per unit length, \( \rho_0 \) is the ambient gas density and \( t \) is the time [53]. Trajectory measurements of a shock can provide a large amount of information about the state of that shock. The Sedov-Taylor blast wave is inherently non-radiative which assumes that the total amount of energy in the blast wave remains constant throughout its evolution [25]. A radiative blast wave is one where radiative cooling of the system occurs on a shorter time scale than the hydrodynamic timescale such that \( \tau_s \sim \Delta R/u_s \). Typically, radiative blast waves are not self similar such that a solution of the form \( R(t) \propto t^\alpha \) cannot be used.
However, there exist two types of radiative blast wave that are self similar, they are known as the Pressure-Driven Snowplow (PDS) and the Momentum-Conserving Snowplow (MCS).

In the PDS case only the thin shell behind the shock, not the interior material, is radiatively cooling. The shock is being driven forward by the high pressure from the non-radiative interior material. In the MCS case both the thin shell behind the shock and the interior material are radiatively cooling, the blast wave sweeps up more mass as it propagates through the ambient medium since the internal material is sufficiently radiative that the internal pressure is negligible [61]. The equations governing the deceleration parameters in a spherical geometry for PDS and MCS are:

$$\alpha = \begin{cases} \frac{2}{\gamma - 1 - k\rho_0} & \text{PDS} \\ \frac{1}{1 - k\rho_0} & \text{MCS} \end{cases}$$  

(2.23)

With $\gamma_i$ the polytropic index at the interior of the blast wave and $k\rho_0$ defined by the density distribution through which the blast wave is propagating [59]. The experiments in this thesis are conducted in a largely uniform medium so that $k\rho_0 = 0$ and as usual the polytopic index is taken to be $\gamma_i = 5/3$. Using these values, we arrive at the Table 2.1 showing the deceleration parameters for spherical and cylindrical geometries in the ST, PDS and MCS cases.

<table>
<thead>
<tr>
<th>Geometry</th>
<th>ST</th>
<th>PDS</th>
<th>MCS</th>
</tr>
</thead>
<tbody>
<tr>
<td>Spherical</td>
<td>2/5</td>
<td>2/7</td>
<td>1/4</td>
</tr>
<tr>
<td>Cylindrical</td>
<td>1/2</td>
<td>3/8</td>
<td>1/3</td>
</tr>
</tbody>
</table>

Table 2.1: Deceleration parameters for spherical and cylindrical blast waves in the Sedov-Taylor, Pressure-Driven Snowplow and Momentum-Conserving Snowplow modes.

2.3.4 Astrophysical Shocks and Scaling from the Laboratory.

In the astrophysical environment there exist many types of shocks and their complex, highly dynamic nature is one of the drivers of evolution on this massive scale. Furthermore, there exists a number of scaling arguments that can be applied to laboratory based shock experiments which can access the same physical mechanisms as those on the astrophysical scale. The cluster blast wave experiments in this thesis
will eventually be evolved into a test bed upon which astrophysical shock physics will be studied. Shocks can be referred to as radiative if the radiative flux is much greater than material energy flux, as is often the case when dealing with cluster blast waves in high Z gases.

A material is considered optically thick if a photon undergoes many interactions while passing through it, whereas an optically thin material will allow a photon to pass through it with very few interactions [25]. When dealing with an optically thick plasma, the simplest case, we can approximate the radiative flux to that of a black body. In such a case the black body radiative flux will scale as $\sim T^4 \sim u_s^8$ as opposed to the weaker temperature scaling of the material energy flux, which scales as $\sim T^{3/2} \sim u_s^3$ [25]. Clearly this change in the energy loss mechanism will lead to a substantial effect on the temperature and density of the upstream material and so the optical thickness of the upstream and downstream material has a great effect on the overall evolution of the system.

Often when discussing shocks it is common to refer to the optical thickness of

![Figure 2.3: Diagram showing the different optical thickness regimes of shocks according to the Drake classification [25]. Some astrophysical and laboratory plasmas are shown in the different regions as examples of the optical thickness’s of such varied phenomena [62, 63]. The reader should also note the evolution of a Supernova from the thick-thick regime to the thin-thin regime of the Supernova Remnant, along the red arrow. Adapted from [6].](image)
downstream and upstream material separately. Standard nomenclature is that a material which is thick downstream and thin upstream is thick-thin. Fig. 2.3 shows a diagram of the different regimes of optical thickness and it also shows where a number of astrophysical shocks and laboratory experiments lie in terms of their optical thickness.

Radiative shocks, particularly those in the thin-thin regime, exhibit extremely complex behaviour which is difficult to model. Modelling is difficult precisely because the radiation will have an effect so far from the shock, the primary radiation source, which then ultimately alters the behaviour of the shock itself as it propagates through that material. Since such radiative shocks are common in the astrophysical domain and are often drivers of its evolution they are of great interest. However, the evolution of astrophysical objects can take hundreds or thousands of years making it extremely difficult to accumulate enough data to quantitatively understand the objects’ evolution. Furthermore, when conducting observations using, for example, a telescope, we cannot select the starting conditions of the system and so rely to a great extent on luck to observe the precise conditions we need to further our understanding. However, in the laboratory we can potentially conduct well characterised and carefully tailored experiments. In order to make such experiments relevant a set of rigorous scaling laws must be applied.

The most commonly used scaling argument employs a number of dimensionless variables which when matched between the laboratory and the astrophysical case mean that the same physical processes dominate in both systems. These early scaling arguments were first developed by Ryutov et al in the late nineties [64, 65]. The Euler Equations, including magneto-hydrodynamic effects, are shown here:

\[
\frac{\partial \rho}{\partial t} + \nabla \cdot (\rho \mathbf{v}) = 0 \tag{2.24}
\]

\[
\rho \left( \frac{\partial \mathbf{v}}{\partial t} + \mathbf{v} \cdot \nabla \mathbf{v} \right) = -\nabla p - \frac{1}{4\pi} \mathbf{B} \times \nabla \times \mathbf{B} \tag{2.25}
\]

\[
\frac{\partial \mathbf{B}}{\partial t} = \nabla \times \mathbf{v} \times \mathbf{B} \tag{2.26}
\]

We complement this with the energy equation:

\[
\frac{\partial \epsilon}{\partial t} \mathbf{v} \cdot \nabla \epsilon = - (\epsilon + p) \nabla \cdot \mathbf{v} \tag{2.27}
\]

Here the equations are written in CGS notation [65]. Where \( \rho \) is the density, \( \mathbf{v} \) is
the velocity, $p$ is the pressure, $\mathbf{B}$ is the magnetic field and $\epsilon$ is the internal energy per unit volume. We define a set of characteristic quantities for a system: $L^*$, $\rho^*$, $p^*$, $v^*$, $\mathbf{B}^*$ the characteristic length, density, pressure velocity and magnetic field scales. These allow us to define the dimensionless variables required for the scaling argument to hold, these are [65]:

$$\hat{r} = \frac{r}{L^*}, \quad \hat{t} = \frac{t}{L^*} \sqrt{\frac{p^*}{\rho^*}}, \quad \hat{\rho} = \frac{\rho}{\rho^*}, \quad \hat{p} = \frac{p}{p^*},$$

$$\hat{v} = \sqrt{\frac{\rho^*}{p^*} \hat{\rho}}, \quad \hat{\mathbf{B}} = \mathbf{B}/\sqrt{p^*}$$

(2.28)

If we work in a polytropic gas then we know that the internal energy density is proportional to the pressure $\epsilon = Cp$. Also the polytropic index of the gas is given by $\gamma = 1 + 1/C$. Using the relationship between internal energy density and pressure with Eq. 2.27 we arrive at:

$$\frac{\partial p}{\partial t} + \mathbf{v} \cdot \nabla p = -\gamma p \nabla \cdot \mathbf{v}$$

(2.29)

To solve Eq. 2.24, 2.25 and 2.26 & 2.29 a set of initial spatial distributions are defined for the density, velocity, pressure and magnetic field as shown here:

$$\rho|_{t=0} = \rho^* f\left(\frac{r}{L^*}\right), \quad p|_{t=0} = p^* g\left(\frac{r}{L^*}\right)$$

$$v|_{t=0} = v^* h\left(\frac{r}{L^*}\right), \quad B|_{t=0} = B^* k\left(\frac{r}{L^*}\right)$$

(2.30)

We then re-cast Eq. 2.24, 2.25, 2.26 and 2.29 in terms of the dimensionless variables defined in the Eq. 2.28. Recasting in this way results in the set of equations [25, 65]:

$$\hat{\rho}|_{t=0} = f(\hat{r}), \quad \hat{p}|_{t=0} = g(\hat{r}),$$

$$\hat{v}|_{t=0} = v^* \sqrt{\frac{\rho^*}{p^*}} h(\hat{r}), \quad \hat{\mathbf{B}}|_{t=0} = \frac{B^*}{\sqrt{p^*}} k(\hat{r}),$$

(2.31)

From these equations it can be seen that the form of the equations is maintained with the dimensionless variables. Therefore, if the form of the initial spatial distributions $f$, $g$, $h$ & $g$ are maintained between the two system, along with the same values for
\[ v^* \sqrt{\rho^*/p^*} = C_1 \quad \& \quad B^* \sqrt{\rho^*} = C_2, \]

where \( C_1 \) and \( C_2 \) are constants, in each system then the two systems are “identical” despite the difference in their real spatial and temporal scales. Stated formally, if these conditions hold then the two systems are “hydrodynamically similar” and will evolve identically from their initial conditions [65].

The solution shown here is for a smooth and non-radiative case which does not match well with often poorly understood astrophysical phenomena or with laboratory experiments. The scaling laws can however be successfully extended to cope with these more complex needs. However particular care must be taken when a radiative laboratory experiment is being designed in order to ensure that the scaling laws can be properly applied in a similar fashion to that shown above.

The reader should also note that a new method has been discovered to successfully scale between the astrophysical case and the laboratory, based on Lie Group Theory rather than dimensionless scaling arguments. This is a formalism of theoretical physics intended to investigate symmetry properties of partial differential equations. The Homothetic Group (HG) is a subgroup of the Lie Groups which can be used to generate scaling arguments. Using this method is possible to select which parameters you wish to match allowing a number of other parameters to be free, which should make it easier to create experiments which can scale properly. However, proper application of such a group theory method requires substantially more understanding of sophisticated mathematics than the dimensionless scaling arguments and so is not broadly used by the experimental community yet. For further details see these papers [66, 67, 68].

2.4 Summary.

This chapter has sought to introduce the basic theoretical underpinnings of the experiments described in this thesis. We began with the definition of a plasma and the basics of how a simple plasma can be modelled. This was then followed with a description of the broad range of different targets used in laser-plasma experiments and the types of experiments carried out using these different types of targets. This led into a description of the interaction of a laser field with a single electron. This forms the basis upon which all laser matter interactions are understood, including the relativistic correction, which become significant at the laser intensities used in the experiments of this thesis. Then the photo-ionisation of atoms was described, ranging from the low to the high intensity regimes.
At this point we move on to the fundamental shock physics concepts required to understand the bulk of the experiments in this thesis. Starting with a basic description of shock and blast wave formation and then the jump conditions present at the shock front. This is followed by the Sedov-Taylor (ST) solution to a blast wave, this analytical solution forms the basis from which the experiments in this thesis are understood. This description is expanded to include how the ST solution was modified to account for the presence of radiation in the system. Finally the scaling laws used to make the broad subject area of Laboratory Astrophysics possible are briefly described for the sake of completeness.
Chapter 3

Diagnosis of Plasmas.

The plasmas investigated in this thesis were primarily diagnosed using optical probing techniques. Time framed or streaked shadowgraphy, dark-field imaging and interferometric images were taken of the plasmas. This chapter will discuss the application of these techniques with reference to the specific plasma geometries present in the experiments described in this thesis and how this affected the analysis and imaging of the system. This will be followed by a description of the operation of streak cameras and x-ray pinhole imaging.
3.1 Optical Imaging of Plasmas.

Plasmas can reflect, transmit, refract, scatter and absorb electromagnetic waves which impinge upon them. Electromagnetic waves can only traverse the regions of plasma with a density below the critical density $n_{\text{crit}}$. In this region the electromagnetic wave will experience dispersion and refraction. These changes to the probe wavefront result in phase changes which are sufficient that the plasma, despite being transparent, can be imaged down to a minimum free electron density of $\sim 10^{16}$ cm$^{-3}$.

3.1.1 Shadowgraphy and Schlieren Imaging.

When light propagates through a medium with a refractive index, $\eta$, the propagation speed of the light is $c_\eta = c/\eta$ where $c$ is the speed of light. If the medium is homogeneous for the whole of the wave front then it will have no impact on the direction of propagation of the light. If the medium has an inhomogeneous refractive index then the path of the light will be angularly deflected, as can be seen in Fig. 3.1 [69].

![Figure 3.1: Schematic showing the modification of a wave front by an object with a non-uniform refractive index gradient. When projected onto a screen this gives a Shadowgraphy image which is proportional to the second spatial derivative of the refractive index [21, 69].](image)

Consider the 2D phase object shown in Fig. 3.1 illuminated by a planar wave, this object has a positive linear gradient in refractive index along the $y$-axis, such
that $\nabla \eta = (0, d\eta / dy)$. If we launch a series of wavelets toward the object and define the path length as $\Delta x = x_2 - x_1$ then the first wavelet emitted at $(x_1, y_0)$ traverses the object in a time $\Delta t = \Delta x \sqrt{\eta / c}$. The other two wavelets are launched at $(x_1, y_0 + \Delta y / 2)$ and $(x_1, y_0 - \Delta y / 2)$. After having travelled for the same time as the first wavelet, these two have travelled a distance of $c\Delta t / \eta_+$ and $c\Delta t / \eta_-$ respectively. The deflection is described by:

$$\theta(y) = \frac{\eta(y) - \eta_+}{\eta_+ - \eta_-} \frac{\Delta x}{\Delta y} \tag{3.1}$$

Where $\eta$ is the refractive index of the medium, $\eta_\pm = \eta(y_0 \pm \Delta y / 2)$ and the other symbols are as shown in Fig. 3.1. Such deflection cannot be observed in plasmas where the electron number density exceeds the critical density, $n \geq n_{\text{crit}}$, because the light will not be able to traverse this region [21, 69]. In the limit $\delta y \to 0$ Eq. 3.1 becomes:

$$\theta_y = \int \frac{\partial \eta}{\partial y} dx, \quad \theta_z = \int \frac{\partial \eta}{\partial z} dx \tag{3.2}$$

As is clear from this equation the angular deflection is dependent on the refractive index of the object. If the deflected light is projected onto a screen, or imaging device, at a distance $L$ from the object then rays which would otherwise have been detected at $(y, z)$ will instead be detected at $(y' = y + L\theta_y, z' = z + L\theta_z)$. If Eq. 3.2 is substituted into $(y', z')$ we obtain [21]:

$$(y', z') = (y + L \frac{\partial}{\partial y} \int \eta dx, \quad z + L \frac{\partial}{\partial z} \int \eta dx) \tag{3.3}$$

We assume a uniform probe intensity of $I_{\text{probe}}$ and a detected intensity of $I_{\text{detector}}$ with a relationship governed by:

$$I_{\text{probe}} \frac{\partial y}{\partial z} = I_{\text{detector}} \frac{\partial y'}{\partial z'} \tag{3.4}$$

This leads to:

$$\frac{I_{\text{probe}}}{I_{\text{detector}}} = 1 + L \left[ \frac{\partial^2}{\partial y^2} + \frac{\partial^2}{\partial z^2} \right] \int \eta dx \tag{3.5}$$

This creates a shadowgraphy image when projected onto a screen where the intensity at any given point is proportional to $\nabla^2 \eta$, the second spatial derivative of the refractive index in the plane [21, 69].

By adding a lens to the basic shadowgraphy set-up and placing an obstruction at
the focal point of the undeflected rays an image can be produced of the gradient of the change in refractive index. This type of imaging is known as Schlieren imaging, from the German ‘Schliere’ meaning streak. Schlieren imaging derives its name from its first use as a method of viewing streak (schliere) like imperfections in glass [69].

Figure 3.2: Schematic of a dark field Schlieren imaging set up. The undeflected rays from a beam that has traversed an inhomogeneous refractive index region are blocked and the deflected rays are imaged. This records the refractive index gradients in the inhomogeneous refractive index region [69].

By bringing the beam to a focus, the relationship between the intensity and the angular deflection is broken. However, the obstruction blocks the deflected rays only allowing the undeflected rays to form an image on the detector surface. The less any given ray is deflected by the object, the less it is blocked, and the higher its intensity on the detector surface. This means that the Schlieren image shows the density gradients of the object through which the rays have passed. It is possible to invert this process by placing an obstruction at the focus of the undeflected rays so that the more the rays have been deflected the higher the intensity of the image, this is known as Dark Field Schlieren. Throughout this thesis it is the dark field Schlieren technique which is used and it will from now on be referred to only as Schlieren imaging [4, 5, 21].
3.1.2 Interferometric Recovery of Electron Density.

Phase Map Retrieval.

Using interferometry it is potentially possible to retrieve the free electron density of the plasma. We consider two electromagnetic waves travelling together in a vacuum, one propagating through an object with some refractive index, the other continues through the vacuum. This introduces a phase shift between the two waves. These two waves can be treated as the two arms of an interferometer and the variation of the fringes in the resulting interferogram will contain information about the refractive index of the medium through which one of the arms has propagated.

![Figure 3.3: Schematic of a Mach-Zehnder interferometer with a plasma in one of the arms. The plasma introduces a phase shift in the wave that passes through it. This phase shift results in a fringe shift from which the electron density of the plasma can be retrieved [4, 70].](image)

The primary source of refractive index in a plasma are the free electrons. The refractive index of the electrons is described as $\eta_e = \sqrt{1 - \frac{n_e}{n_{crit}}}$ which we can approximate as $\eta_e \approx 1 - \frac{1}{2} \frac{n_e}{n_{crit}}$, where $n_e$ and $n_{crit}$ are the electron density and the critical density of the plasma respectively with $n_e \ll n_{crit}$. Therefore extracting the phase shift is equivalent to extracting the electron density of the plasma. The phase shift in an interferogram is described by $\Delta \phi = \frac{2\pi}{\lambda} \int (\Delta \eta) dx$ where $\Delta \eta = (\eta_e - \eta_c) \approx \eta_e - 1$, combining this with the refractive index of the electrons gives us the path integral for the total phase shift of the optical probe as it traverses the plasma shown here [21]:

51
\[ \phi = \frac{2\pi}{\lambda} \int_P (\eta_e - 1)dx \] (3.6)

In Fig. 3.3 a simple Mach-Zehnder interferometer is shown, which can be used to retrieve the phase information subject to certain symmetry constraints. The two electromagnetic waves are described as \( E_1 e^{i\omega t} \) and \( E_2 e^{i(\omega t - \phi)} \) such that arriving at the detector the wave is described by \( (E_1 + E_2)e^{i\phi}e^{i\omega t} \). The detector registers the intensity distribution of the combined electromagnetic waves given by [21]:

\[ |E_{Tot}|^2 = (|E_1|^2 + |E_2|^2) \left( 1 + \frac{2E_1E_2}{E_1^2 + E_2^2} \cos \phi \right) \] (3.7)

It is the \( \cos \phi \) term which gives access to the phase information from the detected signal. The fringe shifts observed in the interferogram are the measure of phase shift across the plasma. The fringes can be mathematically described as [4, 70]:

\[ g(y, z) = a(y, z) + b(y, z)\cos(k_0z + \Phi(y, z)) \] (3.8)

Where \( a(y, z) \) and \( b(y, z) \) are arbitrary scale factors with the fringe period given by \( Z_0 = \frac{2\pi}{k_0} \), for an unshifted fringe \( \Phi(y, z) = 0 \). Taking the Fourier transform of Eq. 3.8 in \( z \) leads to:

\[ G(y, k_z) = A(y, k_z) + \int (c(y, z)e^{-i(k_z - k_0)} + c.c.)dz \] (3.9)

\[ = A(y, k_z) + C(y, k_z - k_0) + C^*(y, k_z + k_0) \] (3.10)

In this equation \( FT[x(y, z)] = X(y, k_z) \) and complex notation has been used such that \( c(y, z) = b(y, z)e^{i\Phi(y, z)} \). \( C \) and \( C^* \) are the side bands of the transform, at \( \pm k_0 \) respectively. A side band is then isolated using a band-pass filter with a soft edge (e.g. a Hann or Gaussian window [71]). The side band is translated back to \( k = 0 \) and an inverse Fourier transform is performed. From this we obtain \( c(y, z) \) and \( \ln|c(y, z)| \) gives the phase between \( -\pi \geq \Phi \geq \pi \). This can be extended into a continuous function by unwrapping the data removing the discontinuities that normally occur at \( \pm \pi \). There exist a number of unwrapping techniques discussed here [72] and here [73], these often use path integrals through the plasma. A phase map retrieved from a null shot can then be subtracted to compensate for aberration on the beam or other background effects, e.g. neutral gas from a cluster target.
Abel Inversion.

As mentioned in Section 3.1.2, Eq. 3.6 is a path integral along a chord across the plasma. The experiments described in this thesis are cylindrically symmetric which affords the opportunity to use Abel Inversion in order to retrieve the electron density itself from the phase map obtained using interferometry. The symmetry of the system allows us to rewrite Eq. 3.6 in cylindrical coordinates giving $(n_e(x, y, z) \rightarrow n_e(\theta, r, z))$. Having rewritten Eq. 3.6 as:

$$ n_e(r, z) = \frac{\lambda n_{\text{crit}}}{\pi^2} \frac{d\phi}{dy} \frac{dy}{\sqrt{y^2 - r^2}} $$

It has the form of the Abel transform and hence we can deduce the radial distribution of the electrons from the chordal measurement [4, 21]. Where $R$ is the maximum extent of the image being analysed, not the maximum extent of the plasma. When taking the actual measurement of the plasma care must be taken that the interferogram extends beyond the plasma filament to an unperturbed region since the Abel Inversion requires $n_e = 0$ at $r = R$ for successful retrieval. This naturally is a widely applicable technique with many other uses in physics beyond phase image analysis.

Figure 3.4: Diagram showing the geometry of Abel Inversion and the accumulation of phase as the probe beam traverses a cylindrically symmetric plasma surrounded by a region of constant refractive index.
3.1.3 Exemplar Images of a Single Blast Wave.

Fig. 3.5 shows examples of Schlieren, Interferometric and Shadography images taken of single blast waves launched using the Blackett Laboratory Laser Consortium (BLLC) Nd:Glass laser system. The laser was operating at $\sim 400 \text{ mJ}$ in 1 ps firing into clustered Argon gas with a backing pressure of 30 Bar using a Parker 99 valve, described in Sec. 4.3.2. They are included in order to give an impression of the appearance of the shocks typical of the experiments described in this thesis. The reader should note that the small round objects seen on the interferometry and shadowgraphy are in fact dust flecks present on the imaging CCD arrays.

Figure 3.5: Exemplar Schlieren (a), interferometric (b) and shadowgraph (c) images of a single blast wave. These blast waves were launched by a vertical heating beam propagating from the top of the image to the bottom with $\sim 400 \text{ mJ}$ of energy on each shot in Argon with a backing pressure of 30 Bar. Here the 527 nm probe beam arrives $10 \pm 0.1 \text{ ns}$ after the heating beam propagated through the medium. The shadowgraph was acquired by blocking the reference arm of the interferometer. The semi-circular shadow visible on the right of the images comes from the gas jet nozzle.

3.1.4 Time Resolved Imaging.

All of the imaging techniques described so far produce 2 dimensional data integrated over the time period of the optical pulse that generated them. In the work described in this thesis the pulse is typically so short that no significant dynamics occur during the integration time. A time history can be constructed by using multiple imaging beams which traverse the plasma at different times creating a series of discrete images. However, the information between the frames is lost. It can sometimes be more useful to acquire continuous temporal information about the evolution of a plasma by ‘streaking’ a 1 dimensional image of the plasma in time. This is achieved
using a streak camera; a schematic of the functional components of a streak camera is shown in Fig. 3.6.

Figure 3.6: Schematic showing the basic working components of a streak camera. A 1D slice is selected from the imaging beam. These selected photons are then converted to electrons at the cathode. These electrons are accelerated and then travel between a pair of plates with a ramping voltage on them creating a ramping field and displacing the electrons further for later times. The electron beam falls on a phosphor screen that converts the signal back to photons which are imaged using a CCD camera with a large chip.

An imaging system is built from the plasma to the streak camera, this can then be used to image self emission or back lit to produce a Schlieren image or shadowgram as described in Sec. 3.1. The object is imaged onto the slit of the streak camera and then the 1D image is re-imaged onto the photo-cathode. The photo-cathode is biased to high voltage and hence, via the photoelectric effect, produces electrons. Using electrostatic lenses the electron beam is collimated and accelerated by an anode. This electron beam then passes between the streak plates, which have a time-ramping voltage applied to them. This results in a time-ramping electric field such that electrons arriving later in time are displaced further than the electrons arriving earlier in time. The electrons then impact a phosphor screen and photons
are emitted and imaged onto a CCD camera. This process results in a 1D image streaked through time on the perpendicular axis.

The rate at which the voltage ramps determines the spread of time in the electrons. Controlling this sweep rate determines the temporal resolution and length of the temporal streak window. The spatial resolution is a product of the slit width and the imaging system magnification and resolution.

### 3.2 X-Ray Diagnostics.

The self-emission of a hot plasma (10’s to 100’s eV) will typically be in the soft x-ray regime and can be imaged using x-ray sensitive cameras or film packs [21, 22]. The detection efficiency of any imaging device typically has a dependence on the wavelength being observed. This dependence often becomes more extreme at the shorter wavelengths, particularly beyond the visible. Therefore, more sophisticated manufacturing techniques must be used to improve sensitivity at short wavelengths. Furthermore, the Quantum Efficiency (QE) of the device needs to be factored in when conducting spectroscopic measurements or, as here, as an imaging device. Also in the low x-ray flux regime it may be possible to do single photon counting and indeed to extract an energy value, with significant error, from the single photon and single pixel interaction with a well calibrated camera [74].

A simple but effective imaging technique when investigating x-ray self-emission is to use a pinhole camera [21]. Here a detector is placed behind a small pinhole, ∼100 μm diameter, such that an image is projected onto the detector from the pinhole. The magnification of the pinhole camera is simply the ratio of the distance of the pinhole to the detector and the distance between the object and pinhole. A schematic of this is shown in Fig. 3.7. The resolution of such a camera is a product of the camera pixel size scaled with magnification and the geometric and diffraction limits of the pinhole limited resolution:

![Image](image.png)

Where $L_{pix}$ is the pixel size, $M$ the magnification, DL the diffraction limit and GL geometric limit of the pinhole. The reader should note that the pinholes used on the XRPC in this work are large, such pinholes are also used in many other experiments [75]. In the case of large pinholes with high magnification it is the geometric limit, GL, which dominates the resolution of the system. In fact the work described in
Figure 3.7: Schematic showing the function of a pinhole camera with the imaged object d from the pinhole and the detector D from the pinhole.

Chapter 7 demonstrates this fact well, here a 100 $\mu$m pinhole was used. In this case the resolution including all terms was 103 $\mu$m and including only the geometric term was 100 $\mu$m, so clearly the geometric limit was dominant. The high magnification of the system means that the effective pixel size is small and so the resolution of the detector is high.

Of course an image obtained in this way is integrated along a chord in the plasma in a similar fashion to back lighting with optical probes, Sec. 3.1. Therefore in a full analysis it must be unwrapped and, in the case of this thesis, mapped to a cylindrically symmetric geometry. For further details on this subject and techniques please see [21]. The absorption of the medium surrounding the emitting material will also alter the observed image and must be taken into account.

The technique can be usefully extended by selecting two filters of similar (often adjacent) atomic number with thickness of each selected such that the transmission spectrum with regard to wavelength is as close to identical as possible, except between that K absorption limits [76]. This creates a spectral notch filter, known as a Ross filter, in the x-ray wavelengths which can be combined with an x-ray pinhole camera, making it possible to capture emission levels in narrow wavelength ranges.
Using a set of narrow Ross filters it is possible to determine the intensity of emission at certain wavelength regions. If we then make an assumptions about the source of the x-rays, for example a Maxwellian temperature distribution, we can make a useful interpretation of the shape of the emission spectrum. The Ross filter signals can then be used to mark out points on the chosen emission spectrum, this in turn can give a good estimate of the temperature of the plasma being observed. This measurement is potentially distorted by the fact that the filter matching cannot be perfect and that there is always some residual signal both at the longer and shorter wavelengths passing through the filter [77].

3.3 Summary.

In this chapter a number of diagnostic techniques are presented which allow us to capture a wide range of information about a plasma. All of the techniques described here are used in Chapters 5, 6 and 7. First we explore the basic optical imaging techniques which measure the refractive index of a plasma and hence its free electron density. These are techniques which can be used to obtain information about the profile of measured shock structures and also about the shape and evolution of a whole shock. Further more when combined with a streak camera, also described here, entire trajectories and other transient phenomena can be observed on a single shot. Finally x-ray diagnostics are briefly explored, in particular x-ray pinhole cameras are described. They are a particularly useful tool when attempting diagnose hot plasmas generated at early time soon after the laser-matter interaction. Finally the underlying methodology for using filter pairs to fix points on an emission spectrum, which can then be used to determine a Maxwellian temperature distribution, is discussed.
Chapter 4

Instrumentation and Experimental Methods.

The focus of this thesis is on high Mach number shocks and blast waves launched by the irradiation of atomic clusters in high intensity laser fields. The aim of this chapter is to introduce the central instruments required to create these high energy density environments. This begins with a description of the design and operation of laser systems and is followed by the generation of atomic clusters in supersonic gas jets flows. Measurements of the performance of the Alameda Applied Science Corporation supersonic jet were conducted by the author and measurements of the Peter Paul valve with large bore nozzle were conducted H.F. Lowe with guidance from the author.
4.1 High Intensity Lasers - An Overview.

Maiman demonstrated the first optical laser in 1960 by flash-lamp pumping a ruby rod lasing at 694.3 nm [78] based on the theoretical work of Einstein [79, 80]. This early relaxation oscillator was improved on substantially by the introduction of the first Q-switching, by McClung and Hellwarth in 1962 [81], and then mode-locking, by Hargrove and Pollack in 1964 [82]. Finally in 1985 Strickland and Mourou demonstrated Chirped Pulse Amplification (CPA) [83] based on the theoretical work of Treacy [84] where pulses are stretched in the time domain, amplified and then re-compressed such that short high energy pulses can be created. CPA techniques paved the way to building high power laser systems while limiting non-linear optical damage.

The work in this thesis is conducted in the high intensity regime ($> 10^{16} \text{ Wcm}^{-2}$), which is attainable with table top systems operating at $> 1 \text{ TW} \ (> 10^{12} \text{ W})$ using short pulses ($< 1 \text{ ps}$) and moderate energies ($< 50 \text{ J}$) with tight focusing ($\sim 50 \mu \text{m}$). Such systems are used to drive experiments ranging from launching cluster blast waves to few cycle high harmonic generation. An alternative route to the high intensity regime requires large facilities capable of delivering kJ of energy on target in ns pulses. These systems require large amounts of space to house long chains of large aperture amplifiers which, due to heat loading, fire infrequently. Such systems are normally used for laser fusion and related experiments, however they cost many millions of pounds to build, run and maintain [85].

4.1.1 Q-Switching and Mode-Locking.

The most basic design of a laser is a cavity with a laser gain medium in it and a pump, for example a pair of mirrors around a flash lamp pumped Nd:Glass rod. In this configuration it is not possible to control the formation or temporal profile of the laser pulse. Once the flash lamps have fired and an inversion has occurred the laser system will typically emit a series of pulses with gradually decreasing energy until the gain medium is depleted, this is a so called relaxation oscillator. The first technique developed to control this process was Q-switching, where the “quality” of the laser cavity is manipulated to control the moment at which amplification starts.

In Q-switching the round trip cavity loses are controlled in such a way that lasing does not occur in the medium until the cavity is Q-switched. A cavity where the total loses are much higher than the total gain is said to be low quality (low Q) and where the total gain is much higher than the total loses in the cavity is said to
be high quality (high Q). Initially the cavity is set to have high losses, low Q, and then after the flash lamp pumps the laser medium there is a delay until the cavity is switched to low loss, high Q [86]. When the system has low Q it prevents full laser action from occurring in the medium by keeping the cavity losses so high that any pulse that begins to form from spontaneous emission is lost within no more than a few cavity round trips. This allows the gain medium to build up an extremely high population inversion. When the desired population inversion is reached the cavity can be switched to high Q, which allows the spontaneous emission to start the process of lasing and a pulse forms. This sudden increase in gain results in a steep energy rise in the pulse and this only comes to an end when the gain of the medium is depleted [81, 86]. Q-switching allows relatively long nanosecond duration pulses to be formed.

A laser cavity is capable of supporting more than a single electromagnetic mode. In a laser that is not mode-locked these longitudinal modes can interact with each other in a number of ways. Interference between different modes can lead to beating in the output, giving strong fluctuations in the pulse properties, or indeed if many modes are available continuous wave (CW) operation can occur [87]. However neither of these scenarios provide pulses suitable for high power lasers, but it is sometimes possible to control the behaviour of these modes in such a way that they start to oscillate in phase. If such coherent behaviour is present the in phase cavity modes will interfere constructively and all other modes will interfere destructively. This results in pulses that not only contain most of the energy present in the laser cavity but also with much shorter pulses durations [86, 88]. This large pulse is partially transmitted at each round trip of the laser cavity producing a uniform pulse train. These pulses are distinct from a free running laser in that they are well defined in their energy, temporal form and they arrive in a uniformly spaced pulse train. Such pulses are necessary to achieve proper amplification using conventional laser amplifiers, particularly because a free running system could produce a single unusually high energy pulse resulting in catastrophic optical damage to the amplifiers down stream. Mode-Locking allows for the creation of very short laser pulses at durations <10 fs [89].

4.1.2 Non-Linear Optics.

The electromagnetic wave interacts with matter via the Lorentz Force which is given by \( \mathbf{F} = q(\mathbf{E} + \mathbf{v} \times \mathbf{B}) \), where \( q \) is the charge, \( \mathbf{E} \) the electric field, \( \mathbf{v} \) the particle
velocity and \( \mathbf{B} \) the magnetic field. We assume that in the interaction between the electromagnetic wave and the atom the effect of the magnetic field is negligible and can be discarded at sub-relativistic intensities. The polarisation density equation can then be written as \( \mathbf{P} = \varepsilon_0 \chi \mathbf{E} \), where \( \chi \) is the electrical susceptibility (typically a tensor). At low intensities the polarisation changes linearly with its excitation, however when entering the high intensity regime the atomic response becomes non-linear with the electric field resulting in [90]:

\[
P(t) = \varepsilon_0 \left[ \chi^{(1)} E(t) + \chi^{(2)} E^2(t) + \chi^{(3)} E^3(t) + \ldots \right]
\]

(4.1)

The magnitude of the \( \chi^{(n)} \) terms is dependent on the properties of the medium through which the laser pulse is moving. For a static, D.C., field there is no dynamic behaviour in Eq. 4.1. However, there is dynamic behaviour for a laser field with an temporally oscillating electric component, \( E(t) \). Therefore for a combination of a D.C. field and oscillating \( E(t) \) field there are potentially non-zero terms in the behaviour of the system at the harmonics of the field (0, \( \omega \), 2\( \omega \), 3\( \omega \), etc). The terms that are of interest in this chapter are \( \chi^{(1)} \), \( \chi^{(2)} \) and \( \chi^{(3)} \) the 1\textsuperscript{st}, 2\textsuperscript{nd} and 3\textsuperscript{rd} orders respectively. Where \( \chi^{(1)} \) leads to the linear refractive index, \( \chi^{(2)} \) leads to 2\textsuperscript{nd} harmonic generation and the Pockel’s effect and \( \chi^{(3)} \) leads to self focusing, Self Phase Modulation (SPM) and 3\textsuperscript{rd} harmonic generation.

The Pockel’s effect occurs when a static electric field is placed across a non-centro symmetric medium, this introduces a fast and slow axis in the medium rotating the polarisation of the light propagating through it. Pockel’s cells are built on this effect and are used together with polarisers to control the direction of propagation through a laser system and the selection of single laser pulses, for example in Q-switching. At high intensities the laser pulse itself can alter the refractive index of a medium, this results from the \( \chi^{(3)} \) term. The refractive index of the medium becomes \( \eta = \eta_0 + \eta_2 I_L \), with \( \eta_0 \) the linear refractive index, \( \eta_2 \) the intensity dependent refractive index and \( I_L \) the laser intensity. The magnitude of \( \eta_2 \) is property of the medium through which the laser pulse is propagating. Spatial variations of the laser pulse intensity can result in self focusing effects and temporal variations result in SPM which can also induce spectral broadening.

Self focusing and self phase modulation can easily become destructive effects in a high power laser system where they can increase the intensity to the point where the medium ionises. For design purposes the Breakup Integral or B-Integral, can be used to estimate whether a system will suffer damage from non-linear effects.
equation governing B-Integral is shown here:

\[
B = \frac{2\pi}{\lambda_L} \int_0^L \eta_2(z) I_L(z) dz \quad (4.2)
\]

Where \( \lambda_L \) is the laser wavelength [91]. It is evaluated throughout the whole systems and if \( B > 3 \) then as a rule of thumb the non-linear effects are at risk of damaging the laser system.

**4.1.3 Optical Parametric Amplification.**

In addition to the effects described in Section 4.1.2, which are used for control in the laser system and the generation of probe pulses, non-linear optical techniques can be used to achieve amplification of laser pulses without the level gain narrowing experienced in conventional laser amplifiers, e.g. Nd:Glass rods, which unavoidably increase minimum pulse duration. This technique is known as Optical Parametric Amplification (OPA).

Optical Parametric Chirped Pulse Amplification (OPCPA) is a combination of the more conventional CPA schemes, see Sec. 4.1.4, with OPAs making it possible to achieve the same pulse energies and durations as pure CPA systems but with much improved contrast [92, 93, 94, 95, 96]. It relies on non-linear optical frequency mixing, which allows one pulse to instantaneously extract energy from another. When carefully performed the pulse that is being amplified will conserve its pulse duration and pulse shape, changing only in its amplitude [90].

Through selection of internal conditions in the non-linear crystal it is possible to either carry out the OPA process or the closely related difference frequency generation (DFG). Fig. 4.1b shows a simple schematic of the atomic process which makes OPA and DFG possible. DFG is used to generate a different wavelength of laser light than that produced by the signal or pump laser, here \( \omega_1 \neq \omega_2 \). Typically in the OPA process the output wavelength is the same as the signal such that \( \omega_1 = \omega_2 \), although other wavelengths can be generated [90].

**4.1.4 Chirped Pulse Amplification.**

In building high power lasers, with short pulse durations and high energy levels, a limiting factor is the damage threshold of the optical material used. The accumulation of B-Integral, c.f. Section 4.1.2, can lead to destructive self focusing. To
overcome this danger the pulse is temporally stretched using dispersive optical elements, then amplified and re-compressed, achieving high powers without damaging the system by controlling the intensity as the pulse propagates through the laser chain. In 1985 Strickland and Mourou made the first practical demonstration of pulse compression [83] with frequency chirped pulses, this technique was based on the theoretical work of Treacy [84] and is known as Chirped Pulse Amplification (CPA). A simple diagram of the CPA concept can be seen in Figure 4.2 illustrating this point.

The technique relies on being able to manipulate the frequency chirp of a pulse. A laser pulse is said to be frequency chirped when a frequency dependent phase, \( \phi(\omega) \), is added to the pulse. In the temporal domain this results in a time delay between the different frequency components of the pulse, stretching the pulse in time. Fig. 4.3 shows a schematic representation of a stretcher and compressor, based on reflective ruled gratings.

To stretch (temporally chirp) the pulse it is passed through a set of optical elements, for example the gratings shown in Fig. 4.3, which introduce a path difference between the different frequency elements of the pulse known as Group Delay
Figure 4.2: Schematic representation of how CPA works in terms of the pulse duration, amplitude and the chirp of the pulse, where the grey scale in the pulses shows the dispersion of the frequencies in the pulse i.e. the chirp.

Figure 4.3: Schematic of a simple grating stretcher and grating compressor, where the red and blue lines represent different frequency components of the laser. This illustrates that one frequency component of the beam must travel farther than another such that the pulse is stretched or compressed. Typically modern stretchers and compressors use curved mirrors rather than lenses since they introduce fewer chromatic aberrations. Modified from [5].

Dispersion [84]. In a typical arrangement positive GDD is added and then after amplification the pulse is passed through another set of optical elements which impart negative GDD. This reverses the initial stretch (positive chirp) and ideally returns the pulse to its original duration, hence leading to much greater powers without damaging the laser system by reducing the accumulated B-integral.

The transmission function of the stretcher and compressor elements is given by
\( t(\omega) = e^{i\phi(\omega)} \) which is added or subtracted respectively. Each frequency component of the laser pulse is treated separately across the full bandwidth. The phase, \( \phi(\omega) \) can be written as a power series, as shown here:

\[
\phi(\omega) = \phi(\omega_L) + (\omega - \omega_L) \frac{d\phi}{d\omega}\bigg|_{\omega_L} + \frac{(\omega - \omega_L)^2}{2!} \frac{d^2\phi}{d\omega^2}\bigg|_{\omega_L} + O \tag{4.3}
\]

Where the second term is the group delay (GD), the third term is the GDD and \( O \) are higher order phase terms. The stretcher element, e.g. Fig. 4.3A, is generally configured to add positive chirp to the pulse altering the pulse duration as \( \Delta \tau = GDD \times \Delta \omega \) [4, 5]. The compressor then adds negative chirp to return the pulse to its original state such that the condition given here is satisfied:

\[
\phi_{\text{stretch}} + \phi_{\text{compress}} = 0 \tag{4.4}
\]

However in a real laser system transmissive optical elements between the stretcher and compressor will add some additional chirp to the laser pulse. For this reason the compressor GDD will need to be adapted to compensate for the material dispersion, this then fulfils the condition given here:

\[
\phi_{\text{stretch}} + \phi_{\text{compress}} + \phi_{\text{material}} = 0 \tag{4.5}
\]

Here \( \phi_{\text{stretch}} \) is the phase imparted by the stretcher, \( \phi_{\text{compress}} \) is the phase imparted by the compressor and \( \phi_{\text{material}} \) is the phase imparted by transmissive optics in the laser system. Further, other process that occur in the laser, for example self phase modulation or gain narrowing, which can alter the higher order terms in Eq. 4.3. These result in temporal aberrations to the pulse such as ringing in the time domain which can very seriously degrade the temporal contrast of a laser system [95, 97].

### 4.2 Laser Systems.

The work presented in this thesis was conducted on three high-power laser systems: the Blackett Laboratory Laser Consortium (BLLC) Nd:Glass laser system, the BLLC Cerberus laser systems and the Central Laser Facility Astra-Gemini laser system. The BLLC Nd:Glass laser system, the first CPA laser constructed outside the USA in the late 1980s, was decommissioned during the progression of the author’s PhD research and as such was the longest continually operating CPA laser in the World. The power amplifiers and compressor gratings of the BLLC Nd:Glass
laser system were used to create a stand alone beam line on the Cerberus laser systems, constructed by the author and his colleagues, dedicated to laser-plasma physics experiments. The Astra-Gemini laser system was used to carry out experiments at higher energy levels than those currently available with the BLLC Cerberus laser system, although upgrades to this system are planned.

4.2.1 The BLLC Nd:Glass Laser System.

The BLLC Nd:Glass laser system was a CPA laser system. The seed pulse was produced with a Time-Bandwidth GLX 200 oscillator, which delivered a $\sim 70 \text{ MHz}$ pulse train with a pulse duration of $\sim 250 \text{ fs}$ and $\sim 1 \text{ nJ}$ of energy. It is made up of a mode-locked cavity with an Nd:Glass gain medium pumped by two 1 W 800 nm laser diodes. The seed pulse is stretched to $\sim 1 \text{ ns}$ and injected into a $\sim 30$ round trip flash lamp pumped Nd:Glass Regenerative Amplifier, a Q-switched high gain linear laser cavity, where the seed pulse sees $\sim 10^6$ gain, reaching 1 mJ. This is then injected into a series of 3 power amplifiers and 3 spatial filters before finally being compressed. Under typical operating conditions the final output of the system is 1 J in 1 ps, achieving 1 TW of power on target. Due to heat loading in the amplifiers the systems fires at full energy of 1 J once per minute and can fire at $\sim 100 \text{ mJ}$ once every ten seconds. A schematic of the laser system is shown in Fig. 4.4.

The BLLC Nd:Glass laser was de-commissioned in February 2011. The power amplifiers from the glass laser have been recycled into a new laser system which is under-construction during this time, referred to as the Cerberus Glass Line (CGL). The new front end of this system uses OPA stages, see Sec. 4.1.3, as the system pre-amplifier, as opposed to the old BLLC Nd:Glass laser which used a regenerative amplifier. The majority of the gain in the old Glass Laser came from the regenerative amplifier, however this introduced a significant amount of pre-pulse and gain narrowing. The slicer Pockel’s cell shown in Fig. 4.4 was used to eliminate pre-pulse on a longer few ns time scale, however due to electronic jitter the short time scale pre-pulse ($> 3 \text{ ns}$) cannot be reliable addressed in this manner. Due to the physical process which makes amplification possible in OPAs they can produce far shorter pre-pulses durations at 10’s ps if a sufficiently short duration pump pulse is employed [98].

This new design shifts the majority of the gain in CGL into the OPA stages, which are capable of supporting a large bandwidth with an improved temporal con-
The expected characteristics of this laser system are 10 J and \( \sim 500 \) fs with a high spatial and temporal quality.

### 4.2.2 Astra-Gemini Laser System.

The Astra-Gemini laser system is a Ti:Saphire based CPA laser with twin beam lines seeded from a single master oscillator power amplifier scheme. The Gemini project is an extension of the original Astra laser which delivers 1 J in 40 fs at 10 Hz with a wavelength of \( \sim 800 \) nm. Whilst the Astra laser is a very useful system in itself the experiments carried out at the Central Laser Facility in this thesis revolve around the application of higher drive energies to create stronger shocks, which scale with energy as \( E^{1/4} \) c.f. Sec. 2.3.1. For this reason the full Astra-Gemini system was used, operating at 15 J in \( \sim 40 \) fs at a wavelength of \( \sim 800 \) nm [99]. In order to achieve the higher laser power the shot rate of the system is significantly reduced to...
a shot per minute. It should however be noted that this remains a comparatively high shot rate for shock experiments.

4.3 Cluster Sources.

4.3.1 Generation of Atomic Clusters.

Atomic clusters are a unique target medium for high power laser experiments, which gives both high laser energy absorption and laser penetration into the medium, due to their global gas density and local solid density. Such that when a high intensity laser is focused into the medium it creates a hot plasma filament which can evolve into a cylindrical shock structure. When a gas of atoms, or indeed molecules, expands into a vacuum the small inter atomic forces can become dominant and the atoms can form aggregates known as clusters. The selection of atomic species, gas jet backing pressure, gas jet nozzle geometry and temperature can be used to control the size of the clusters and the percentage of the atomic gas which forms into clusters [100].
The first study of clusters was carried out by Hagena and Obert in 1972 [100] where the size distribution and clustering degree under different pressures, temperatures and nozzle geometries was investigated. This lead to an empirical scaling law:

\[ \Gamma^* = k \frac{\left(\frac{d}{\tan \alpha}\right)^{0.85} P_0}{T_0^{2.29}} \]  

(4.6)

This is the so called Hagena parameter where \( d[\mu m] \) is the nozzle diameter, \( \alpha[\circ] \) is the expansion half angle, \( P_0[\text{mBar}] \) is the backing pressure, \( T_0 \) is the gas temperature pre-expansion in Kelvin and \( k \) is the condensation factor. The condensation factor, \( k \), is determined by the gas species used and is dependent on the bond formation in this given gas [100]. Table 4.1 shows the condensation factor for a number of atomic and molecular gases [101].

<table>
<thead>
<tr>
<th>Gas Species</th>
<th>Condensation Factor, ( k )</th>
</tr>
</thead>
<tbody>
<tr>
<td>( H_2 )</td>
<td>184</td>
</tr>
<tr>
<td>He</td>
<td>3.85</td>
</tr>
<tr>
<td>( N_2 )</td>
<td>528</td>
</tr>
<tr>
<td>( CH_4 )</td>
<td>2360</td>
</tr>
<tr>
<td>Ar</td>
<td>1650</td>
</tr>
<tr>
<td>Kr</td>
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<tr>
<td>Xe</td>
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</tbody>
</table>

Table 4.1: Table showing the condensation parameter, \( k \), for a variety of atomic and molecular gases for determination of the Hagena parameter, Equation 4.6. He is included as a reference gas which does not cluster well and is unlikely to be used as a target gas on its own [101].

The Hagena parameter, Equation 4.6, can be used to estimate the mean size of clusters generated by a gas jet based on the backing pressure and temperature selected. However it is an empirical law determined using small aperture, < 1 mm, gas jets and must be modified when applied to larger aperture gas jets. Clustering normally sets in when \( \Gamma^* > 100 \), the number of atoms or molecules per cluster, \( N_c \), can then be estimated using [47]:

\[ N_c \approx 25 \left( \frac{\Gamma^*}{1000} \right)^{2.40} \]  

(4.7)

As is clear from Eq. 4.6, and by extension Eq. 4.7, the degree of clustering and cluster size varies strongly with both the backing pressure and the temperature of the gas before expansion. Varying the backing pressure then allows control of the cluster size but also varies the gas density of the global gas. However with certain designs of gas jet, e.g. the Parker-Hannifin 99 valve [101], it is also possible to control the pre-expansion temperature of the gas. This makes it possible to vary
cluster size while not greatly reducing the global gas density, Hagena scales as $T^{2.29}$, but also makes it possible to cluster gases with a low k, like hydrogen, expanding the breadth of experimental conditions that can be investigated in clustered gases.

### 4.3.2 Gas Jets Used for Cluster Generation.

A number of supersonic gas jets can be used to generate atomic clusters. These different jets provide access to a cluster medium with varying characteristics. In particular the gas volume and mean cluster size can be controlled in this way but also the pre-expansion temperature and repetition rate are also determined by jet selection. In this section the supersonic gas jets used in the experiments in this thesis are detailed. The jets described are Parker-Hannifin 99 valves (referred to as the Parker 99 valve), Alameda Applied Science Corp. (AASC) Supersonic Large Bore jets and Peter-Paul EH22G7DCCM valves with large custom nozzles.

**Modified Parker 99 valve.**

The experiments described in Chapter 5 make use of the Modified Parker 99 valve. The Parker 99 valve is distinct from the many valves because it uses a poppet made with a diamond compound or composite and a solenoid to make the seal at the nozzle, which makes it possible to cool the Parker 99 with liquid nitrogen down to $\sim 125$ Kelvin using a cooling jacket around the jet body. At these temperatures clusters will form in lower k gases, see Sec. 4.3.1, including Hydrogen.

Fig. 4.6 shows a schematic of the modified Parker 99 valve used in this thesis. The modified nozzle has the same diameter as the standard Parker 99 valve at 500 $\mu$m. The expansion half angle of the modified nozzle is 45°. The Parker 99 is driven by a Parker-Hannifin Iota One power supply producing 25 Volts and set to an opening time of 4 ms.

The expansion of the gas from the nozzle can be modelled as a free isentropic ideal gas expansion as described by Miller [102]. This allows for the calculation of both the temperature, Eq. 4.9, and the average density, Eq. 4.10, of the centre line gas flow. First the Mach number of the gas flow must be calculated according to:

$$M = A \left( \frac{x}{d} - B \right)^{\gamma^{-1}} - \frac{\gamma + 1}{2A(\gamma - 1)(\frac{\gamma}{2} - B)^{\gamma^{-1}}}$$

(4.8)

Where $A$ and $B$ are near constant, varying only weakly with $\gamma$, as $A = 3.26$ and $B = 0.075$ for $\gamma = 5/3$ and $x$ is the distance from the nozzle and $d$ is nozzle diameter.
Using the Mach number the temperature:

\[ T[K] = T_0 \left( 1 + \frac{\gamma - 1}{2} M^2 \right)^{-1} \] (4.9)

And the average density:

\[ n = n_0 \left( 1 + \frac{\gamma - 1}{2} M^2 \right)^{\frac{1}{\gamma - 1}} \] (4.10)

Of the centre line gas can then be calculated. \( T_0 \) and \( n_0 \) are the pre-expansion temperature and density respectively. However it should be noted that when the condition \( x/d > 2.5 \) is not fulfilled then these equations are no longer valid.

Typical operating conditions for the modified Parker 99 valve would be \( \sim 40 \) bar of argon at room temperature (\( \sim 300 \) K). This would, at a distance of 2.5 mm from the nozzle, result in an average density at the centre line of \( \sim 10^{19} \) cm\(^{-3} \) and a temperature at the centre line of 129 Kelvin. Furthermore with the application of the Hagena parameter it is estimated that the mean cluster radius would be \( \sim 35 \) nm and local density within such a cluster is \( \sim 10^{22} \) cm\(^{-3} \).
In previous experiments it has been observed that the fast shock generated in clustered gases using small orifice gas jets, for example the Parker 99 valve, are too small such that the shock propagates to the edge gas volume before the instabilities we are looking can develop. The solution to this is to use large bore gas jets with an orifice $\sim 10$ mm. However the author is not aware of any previous studies available in the literature that investigate the cluster properties of large bore gas jets. In November 2011 the author was loaned a supersonic large bore gas jet produced by the Alameda Applied Science Corporation (AASC).

In order to ascertain the viability of the AASC jet as a source of cluster for experimental use two tests were applied. The Rayleigh scattering of the AASC jet was measured continuously in time over a variety of positions to determine the time to saturation of cluster formation and a pulsed Nd:YAG interferometer was used to determine the gas density profile of the jet at saturation. Previous experience informs us that the valve must open for a reasonable long time for cluster formation to peak and stabilise, for this reason two return springs were tested to optimise performance.

To retrieve the gas density profile a Mach-Zehnder interferometer was used, this setup is shown in Fig. 4.8A. A Coherent Powerlite Nd:YAG laser was used with a pulse duration $\sim 9$ ns doubled to a wavelength of 532 nm. The YAG arrival was timed to correspond to the peak of cluster formation, the time at which experiments are launched in the clustered gas. The gas density can be retrieved from an interferometric image in the same manner as the free electron density in a plasma is, this process is described in Sec. 3.1.2, where the phase map imprinted on the interferometric fringes is Abel inverted assuming cylindrical symmetry along the gas jet axis.

The reader should note that the clusters present in the gas are likely to change the refractive index of the medium which will alter the phase map. The author is aware of no comprehensive study which quantifies this effect however the necessary correction will be small and falls within the error of these measurements.

Fig. 4.7 shows density profiles retrieved by the AASC team for the jet tested for its cluster properties by the author. The densities retrieved over a range of pressures by the author do not differ significantly from those measured by AASC. As is clear from Fig. 4.7 the jet forms a near uniform cylinder of gas which falls off slowly with distance from the nozzle. The uniformity of this volume is a great asset in terms of
modelling experiments since the density gradient is low. Furthermore the densities of $10^{18}$ to $10^{19}$ cm$^{-3}$ are still experimentally useful for the blast wave experiments described in this thesis, even if they are somewhat lower than those produced by the Parker 99 valve close to the nozzle, see Sec. 4.3.2.

Using the Hagena parameter, Eq. 4.6 we can estimate the cluster size from the nozzle. With Argon at room temperature with a backing pressure of 50 bar we would estimate the Hagena parameter to be $\Gamma \sim 106$ giving a mean cluster size of $N \sim 100$ atoms per cluster or radius $\sim 150$ nm. Since the Hagena parameter is an empirical formula that was derived using small gas jet nozzles it is unlikely to apply well here and will probable overestimate the cluster size. It does however indicate that we would expect strong clustering of the gas from the AASC jet.

Rayleigh Scattering (RS) occurs when photons interact with an object much smaller than their wavelength leading to an elastic scattering event. The intensity of the scattered light is related to the size of the object from which the photons scattered. Therefore the signal produced by the RS for clusters scales non-linearly with the size of the clusters and the percentage (degree) of the monoatomic gas that
Figure 4.8: Schematic of the (A) gas density and (B) clustering profile retrieval set-ups for the Alameda Applied Science Corporation supersonic jet.

has aggregated into clusters [52, 103]. The equation governing the cross section of RS is:

$$\sigma = \frac{8\pi r^6}{3\lambda^4} \left( \frac{n^2 - 1}{n^2 + 2} \right)^2$$  \hspace{1cm} (4.11)

Where \( r \) is the radius of the object, \( n \) is the number density and \( \lambda \) is the wavelength of the probe. As is clear from this equation it is the wavelength of the probe which most strongly determines whether an object can be observed or not since neither the density nor the cluster radius is an independent variable.

The AASC jet testing was performed in a vacuum chamber at \(< 4 \times 10^{-4} \) mBar. A Sumix M72 CMOS camera was used to visualise the spatial extent of the scattering signal, and hence the cluster volume, and the temporal evolution of the RS signal was monitored with a ThorLabs DET10A photodiode. This diagnostic arrangement is shown in Fig. 4.8B. The valve was tested using two spring stiffness’s 1170 lb/in and 585 lb/in driven at a supply voltage of 530 V or 720 V, over a range of backing pressure and positions relative to the nozzle. The RS probe beam was produced by a diode pumped doubled Nd:YAG at 532 nm with a power \( \sim 300 \) mW which was focused through the gas volume with a cylindrical lens. The use of a loose line focus makes it possible to determine the cluster signal with respect to the distance from
the nozzle.

Changing the spring stiffness affects the return of the magnetically driven flyer plate and hence the opening time of the jet. The AASC jet was designed to have a very short opening time in its standard operating conditions. A softer spring increases this opening time which is better suited to steady state cluster formation.

Figure 4.9: Temporal evolution of the Rayleigh scattering signal of AASC jet with 30 bars of argon with the reservoir at room temperature. Measured on a Thorlabs DET10 photodiode and a line focused 532 nm probe beam. The red line used the soft spring and a supply voltage 720 V, the green line the soft spring with 530 V and the stiff spring with 530 V.

The RS signal from the AASC jet under different operating conditions is plotted in Fig. 4.9, the reader should note that the signal strength in this plot is arbitrary and should not be compared between the different operating conditions. When used in an experiment the laser pulse will usually be timed to arrive when the cluster formation saturates because this leads to greater shot to shot stability. The saturation of clustering leads to a plateau in the RS signal which can be seen in Fig. 4.9 for the trace corresponding to the soft spring (585 lb/in) driven at 720 V. From the RS scattering it is clear that a higher than standard drive voltage is not required to reach the plateau however the use of the stiff spring (1170 lb/in) results in the valve closing too quickly and inhibiting saturation of the clustering process. Therefore from this point onwards and in the experiments described in Chapters 6
and 7 the 585 lb/in spring is used. Whilst it is not experimentally important it is of note that the saturation region last for \(\sim 3\) ms with a drive voltage of 530 V and \(\sim 4\) ms with a drive voltage of 720 V, which is slightly shorter than the Parker 99 valves saturation time of \(\sim 4\) ms.

Figure 4.10: Graph showing the Rayleigh scattering signal as compared to backing pressure for the soft spring. The soft spring measurement was taken 4 mm from the nozzle. The error bars cannot be seen on this graph because they are too small.

Figure 4.11: Graph showing the Rayleigh scattering signal as compared to backing pressure for the stiff spring. The measurement was taken 1.4 mm from the nozzle. The error bars cannot be seen on this graph because they are too small.
Hagena’s work tells us that the size of the clusters grows with increased backing pressure which results in a RS signal which increases as a power law with respect to the backing pressure [100]. The RS signal scaling according to Hagena goes as $I_{RS} \sim P_0^{3.35}$ where $I_{RS}$ is the RS intensity and $P_0$ is the backing pressure. However the much larger nozzle diameter of the AASC jet will alter the power law coefficient as can be seen in Fig. 4.10 and 4.11 where the RS signal with respect to the backing pressure for the stiff and soft springs are plotted respectively.

![Table 4.2: Table showing the power law coefficients of the RS versus backing pressure at different positions for the AASC jet with the soft spring driven at 530 V.](image)

<table>
<thead>
<tr>
<th>Distance from nozzle / mm</th>
<th>Power Law Coefficient</th>
</tr>
</thead>
<tbody>
<tr>
<td>4</td>
<td>2.23</td>
</tr>
<tr>
<td>6</td>
<td>2.43</td>
</tr>
<tr>
<td>8</td>
<td>2.52</td>
</tr>
<tr>
<td>10</td>
<td>2.51</td>
</tr>
<tr>
<td>12</td>
<td>2.66</td>
</tr>
</tbody>
</table>

The power law that governs the growth of RS signal with respect to backing pressure grows significantly more slowly than with small diameter nozzles. However it is also clear from Figures 4.10 and 4.11 that the onset of clustering begins at very small backing pressure, with signal observed with as little as 2 bar backing the jet. The author is not aware that clustering has ever been observed with so little backing pressure in the past. Given in Tab. 4.2 the soft spring power law coefficients over range of distances from the nozzle.

It is possible to approximately determine the mean size of the clusters produced by the AASC jet. If we assume that scattering signal is only observable on the photodiode when the clusters contain $>200$ argon atoms, a limit set by the wavelength of the laser as shown in Eq. 4.11. Further we assume the clusters to scale in a similar fashion to that predicted by Hagena, where $N = \text{Const} \times P_0^{2.35} = \text{Const} \times P_0^{RS-1}$, and that substituting in the measured coefficient we can approximate the mean cluster size [52].

Choosing to consider the power law at an experimentally useful 4 mm from the jet nozzle, see Tab. 4.2, and subtracting one from the power law coefficient we can determine that constant $= 85.1$. Therefore the mean cluster size scales as $n = 85.1 \times P_0^{1.23}$ with the backing pressure. Fig. 4.12 shows the graph of the
Figure 4.12: Graph showing the estimated mean clusters size with respect to the backing pressure 4 mm from the AASC jet nozzle with the 585 lb/in spring driven at 530 Volts. Here it is assumed that the scatter signal is only visible when clusters contain >200 Argon atoms.

estimated mean clusters size with respect to the backing 4 mm from the nozzle with the 585 lb/in spring driven at 530 Volts. As can be seen in this figure we would predict a mean cluster of size of >10,000 argon atoms per cluster at a backing pressure of 50 bar compared to only 12 atoms per cluster with the Parker 99 valve.

Figure 4.13: a) Composite image showing the spatial extent of the RS signal, and hence the cluster volume, taken at 4 mm, 9 mm & 12 mm from the AASC jet nozzle. b) Graph showing a line-out taken from the RS signal measured on the camera 4 mm from the nozzle. It is clear from this that the clusters occupy a spatial volume with a diameter of ~10 mm. All images were collected in Argon with a backing pressure of 50, using the 585 lb/in spring driven at 530 Volts.
In addition to measuring the RS signal on a photodiode a camera was also used to image the spatial distribution of the clusters in the gas jet. The shutter time on the camera is relatively long (10's ms) compared to the jet opening time so the observed signal is temporally integrated. Fig. 4.13a shows a composite image constructed from shots at 4 mm, 9 mm & 12 mm below the gas jet nozzle. All images were collected in Argon with a backing pressure of 50 bar, using the 585 lb/in spring driven at 530 volts. Fig. 4.13b shows a line-out taken through the RS signal 4 mm from the nozzle. This line-out shows that 4 mm from the nozzle the clusters occupy a near uniform volume with a diameter of $\sim 10$ mm. Line-outs taken from the RS signal further from the nozzle show a slow expansion of the cluster volume as a function of distance from the jet.

The gas volume diameter measured by the interferometry, Fig. 4.7, is slightly larger than that of the RS signal. Furthermore the expansion of the gas and cluster volumes follow the same trend. This suggests that the clusters have become ‘frozen in the flow’ of the gas forming a cylinder of clusters. This leaves a large uniform cluster volume to deposit energy into and also a large gas volume for the resulting shock to propagate through.

**Peter Paul valve.**

Much of the Astra-Gemini campaign was carried out using a Peter Paul EH22G7 DCCM valve mounted with a large bore (10 mm aperture) custom nozzle. Testing of the clustering properties of this valve and nozzle is still on going. However from the extremely strong laser energy absorption properties, measured in Sec. 7.1.1, it is clear that there must be a large number of clusters present in the gas.

It is of great importance to measure the neutral gas density of the medium in order to understand the propagation dynamics of shocks launched from the cluster plasma and indeed the behaviour of the early time plasma generated by the laser. This measurement was taken using the same interferometric method as that described in Sec. 4.3.2.

Shown in Fig.4.14 is the measured radial profile of the neutral gas density of the Peter Paul valve with large bore nozzle. This line-out was taken 1 cm from the end of the nozzle in argon with a backing pressure of 40 bar. This distance from the nozzle is typical of the experimental arrangement used in Chapters 6 and 7. The level of noise observed in this image is a product of the relatively low density, below $10^{17}$ cm$^{-2}$ interferometric density retrieval becomes problematic due to the small
fringe shift.

The measured neutral gas density profiles of the Peter Paul valve, Fig. 4.14, share a number of characteristics with the AASC jet’s profile, Fig. 4.7. Most notable the profile has steep edges indicating, again, that we have a well formed cylinder of gas and a slow density gradient with distance from the nozzle tip. This is ideal for blast wave experiments since it prevents anomalous shock velocity progression as the spatial density change affects the sound speed of the medium. Also a slow gradient was measured in the fall off of density with distance from the nozzle.

![Graph showing the measured neutral density radial profile](image)

**Figure 4.14:** Graph showing the measured neutral density radial profile produced in argon at 40 bar with the Peter Paul valve and large bore nozzle 1 cm from the end of the nozzle.

Discussions with S.P. Mangles indicate that Peter Paul valves typically have a very linear scaling between backing pressure and neutral density. Therefore the measured neutral density at 40 bar can be used along with a number density of 0 cm\(^{-3}\) at 0 bar of backing pressure to create the graph shown in Fig. 4.15 where an estimate of the neutral density at a given backing pressure is presented. This line shown here obeys the equation \(n = P \ast 6 \times 10^{15}\), where \(n\) is the neutral gas density in cm\(^{-3}\) and \(P\) is the backing pressure in bar. Characterisation of the Peter Paul valve with the large bore nozzle to confirm this relationship is on going.

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4.4 Summary.

Presented in this chapter is a brief description of the primary technologies underlying the laser-cluster interaction experiments conducted in this thesis. Firstly the basic techniques used to generate short high power laser pulses are discussed, ranging from the building blocks of simple optical amplifier techniques to the chirped pulse amplification to the more recent development of the optical parametric amplifier combined with the CPA technique to form OPCPA. Focus is given to the damage threshold conditions of high power laser systems and also a discussion of the demands for high fidelity re-compression of laser pulses in the CPA scheme is discussed. This is then joined by a description of the key features of the laser systems used to conduct the experiments described in this thesis and a schematic of the Glass beam line of the IC Cerberus laser system, which is currently under development.

The target medium deployed on all the experiments described in this thesis are atomic cluster gases. For this reason an overview of the formation of atomic clusters is given include the empirical Hagena scale law. Then the actual gas jets used to produce the atomic clusters in relevant experiments are described. Particular focus is given to the newly available AASC supersonic jet which was characterised by the author and found to be an exciting new source of large atomic clusters which also provides a large near uniform target volume for experiments. Characterisation of the Peter Paul with large bore nozzle, used on the experiments described in Chapters 6 and 7, is also explored.
Chapter 5

Characterisation Of Upstream Medium.

In this chapter we discuss the first reported experimental study of the cluster content of the upstream medium from the laser heated clusters. This is a study of interest to work presented in this thesis, and more broadly to those who use clusters as a target, since the presence of clusters upstream of the laser heated plasma could potentially alter the radiative properties of the blast waves. This work was carried out during mid 2010 on the Blackett Laboratory Laser Consortium Nd:Glass laser. The experimental work was carried out by the author and H.W. Doyle and the data analysis was carried out jointly by H.W. Doyle and author. It took advantage of a perpendicular heating beam geometry implemented for the study of proton probing of shocks described in reference [6]. This work has been published by the author in High Energy Density Physics [57].
5.1 Motivation for the Study of the Upstream Medium.

It has been suggested that the blast waves launched in clustered media could be used to conduct laboratory astrophysics studies on radiative shocks. Since initially the gas medium is clustered and clusters will, if not interfered with, survive for a much longer time than the blast wave propagation time we must investigate the extent to which clusters are present in the upstream medium.

It was feared that the opacity of any significant number of clusters in the upstream medium would alter the radiation dynamics of the system to the extent that such experiments will no longer constitute a well enough understood system to create well scaled laboratory astrophysics experiments. In particular an increase in absorption of the upstream medium would result in an increased precursor temperature and hence a faster shock propagation speed.

The author is not aware of any previous study explicitly aimed at investigating the medium upstream of a shock launched into a cluster medium. However the work of A.S. Moore et al. [104], which describes the radiation emitted by a propagating blast wave launched in a clustered medium, led us to previously believe that the clusters were being disassembled by the emitted radiation immediately upstream of the shock. Although this was never experimentally confirmed nor is there theoretically calculated basis for this assertion.

However if the radiation from the shock is being absorbed by the upstream clusters then the optical depth would potentially shorten from that of neutral or lightly ionised gases of the same species with the same number density and temperature. Additionally previous work using multiple beams in clustered media appears to be limited to D.R. Symes’ “machined” cluster medium work [105] and J. Lazarus’ near parallel colliding blast wave experiments [4]. In Symes’ experiments a long pulse laser was used to create an interference pattern which delivered enough energy to disassemble the clusters at the peaks of the field leaving behind “slices” of clustered and un-clustered gas when the second, short pulse beam, arrived [105]. This resulted in a series of near spherical blast waves which collided with each other. In Lazarus’ work a Fresnel bi-prism was used to split a single short pulse creating, two near parallel blast waves which collided with each other [4]. This collision resulted in a density enhancement at the collision point, a similar feature is observed in the experiments described in this chapter.

In this experiment two perpendicular beams are focused into the cluster medium with a variable time delay between the beams. It was found that as the delay between
beams was increased the region into which the second beam could deposit energy would change. A longer delay leads to a larger gap between the first beam and the deposition region of the second beam. The hallmark of a clustered medium is strong laser energy absorption and so the experiment indicates that the clusters around the first beam are being disassembled such that the second beam cannot deposit energy near the first beam. This region grows over time at a speed which can be measured and leads to the conclusion that fast ions rather than hydrodynamic flow, radiation or electrons are responsible for the change in the upstream medium.

5.2 Experimental Set-Up and Method.

The experiments described in this chapter were carried out at Imperial College using the Blackett Laboratory Laser Consortium Nd:Glass laser system. The system was run at a shot per minute with ~1 J in 1 ps at a wavelength of 1054 nm. The output beam from the laser was split 50/50 to produce two ~500 mJ beams with a variable delay between the heating beams and delivered an on target intensity of ~10^{17} \text{ Wcm}^{-2}. All of the laser pulses in the experiment are derived from a single pulse of the laser oscillator. This means all beams are optically synchronised preventing electronic jitter from interfering with the accuracy of the measurements.

A small amount of the main beam is split off, using a 4% wedge in the main beam, and then focusing through a second harmonic crystal (KDP) producing a short, <1 ps, 527 nm probe pulse which is then split into the two linear polarisation states, with one delayed with respect to the other. The delayed pulse is re-injected into the probe line and through the interaction region and imaging system before being split by a polariser onto separate CCD cameras. This makes two time frames available for viewing on a single shot.

The heating beam was injected through the top of the experimental chamber and the second heating beam was injected through the front of the chamber with a variable delay between them. This means that the beams cross paths at the centre of the chamber, where they are both brought to a focus, perpendicular to each other forming a cross. The optical probe is injected through the side of the chamber and so is perpendicular to both of the heating beams. Spatial overlap was achieved by obscuring the three beams on a 10 \mu m wire. A schematic representation of the experimental set-up is shown in Fig. 5.1.

Both of the heating beams are focused using f/3 GRIN lenses giving a focal length of ~13 cm. A GRIN lens uses an axially varying refractive index to eliminate
spherical aberration found in typical singlet lenses and hence improves the focus. The probe beams are collimated through the interaction region and then collected using a large \( f/4 \) lens to capture as many image rays as possible. The imaging system had a \( \times 2 \) magnification and before the final imaging optic the beam was split using a 50/50 beam splitter. One arm of this imaging system was used to capture an interferogram and the other was used to capture a dark field Schlieren image of the interaction. As is clear from Fig. 5.1 the Mach-Zehnder interferometer was used in order to maximise the field of view. In the Dark-Field Schlieren arm either a wire of \( \sim 1 \) mm diameter or an unravelled paper-clip were used as the stop. For details of the use of dark field Schlieren and interferometric imaging please see Sec. 3.1. A Newport Navy resolution slide was used to calibrate the spatial resolution of the imaging systems. Additionally the reader should note at this point the experiment is not cylindrically symmetric and as such interferometry can only be used to determine the electron density away from the region where the two heating beams cross breaking the symmetry, further more the fringes were always set horizontally in this experiment so only the electron density of the vertical shock could be measured in this geometry.

The vertical heating beam is used as the \( t_0 \) reference in this experiment. By varying the delay on the green probe beam it is possible to bracket the fast initial ionisation event caused by the interaction of the clusters and the vertical heating beam. The point at which the ionisation is only just seen to appear on the imaging CCD cameras is defined as \( t_0 \). To time in the horizontal heating beam we observe the cluster breakup effect being studied in this experiment. If the two beams are not temporally overlapped then the cluster breakup effect leaves a gap between the two blast waves, so at temporal overlap there is no gap between the two blast waves creating a cross pattern. From here the horizontal beam can be delayed in time relative to each other to investigate the rate at which this gap grows.

A modified Parker 99 valve was used, see Sec. 4.3.2, with argon gas. Argon is a relatively inexpensive gas which clusters easily at room temperature and at modest backing pressure, see Sec. 4.3.1, with a reasonably high atomic number of 18. The opening time of the gas jet was set to 4 ms with heating and probe beams arriving \( \sim 3 \) ms after the opening of the valve in order to optimise cluster formation in gas in the less turbulent late time flow. The backing pressure was set to 30 Bar in all experiments which results in an atomic density of \( \sim 10^{18} \text{ cm}^{-3} \) and a mean cluster size of \( \sim 10 \) nm.
Figure 5.1: Schematic representation of the set-up of the perpendicular blast wave experiment showing both the heating and probe beam geometries. The heating beams have a pulse duration of 1 ps with an energy of $\sim 1$ J at a wavelength of 1054 nm. The two heating beams are both brought to focus at the centre of the chamber perpendicular to each other. The 527 nm optically synchronised probe beam, pulse duration of $< 1$ ps, is split by polarisation into two beams $s$ and $p$. A delay is introduced between these two pulse such that two time frames are captured on each shot, not shown in the figure for simplicity.

When a time delay is present between the two heating beams a gap forms. This is associated with the disassembly of clusters in the medium and the reduction in laser energy absorption for the 2nd beam. To study this effect the time delay between the two beams was scanned, the error in the heating beam delay steps was $\pm 3$ ps. Multiple shots at each heating delay step were used to acquire the average gap size. The fluctuation in gap size was primarily a result of the shot to shot energy variation originating in the regenerative amplifier of the laser system. Further more the probe delay was also scanned at each heating delay step.

5.3 Results, Analysis and Conclusions.

In order to form a shock in this system we rely on the efficient absorption of laser energy, $>70\%$, by any atomic clusters present in the volume through which the heating laser propagates. Therefore a lack of clusters will result in a substantially reduced energy absorption and no shock will form. Using this basic premise we can
observe that there is a region around the vertical, early time, blast wave which does not contain a significant number of clusters, hence inhibiting the creation of the horizontal, late time, blast wave.

Figure 5.2: Interferometry (a-c) and Schlieren (d-f) images of dual blast waves in the perpendicular blast wave geometry imaged 10±0.1 ns after the arrival of the vertical heating beam. Here the horizontal heating beam is delayed with respect to the vertical heating beam with the time delay increasing from left to right. In (a) the horizontal beam arrives before the vertical by ∼30 ps. In (b) the horizontal and vertical beams arrive simultaneously. In (c) the horizontal beam is delayed by 3 ns. The delay in (d) is 10 ps, in (e) 2.3 ns and in (f) 4.2 ns. Since free electron density retrieval using a single interferometry beam relies on cylindrical symmetry the density at the meeting points cannot be determined here. However the Schlieren imaging can be used to resolve the blast wave shell position accurately. Due to the density gradient from the gas jet the two sides of the blast wave do not behave identically however with increased delay a growth in the gap between the horizontal and vertical blast waves can be observed on both sides.

By plotting the growth in the gap size with increased time delay we can determine the velocity of the disassembly wave causing the break up of clusters in front of the vertical blast wave. Fig 5.4 shows the plot of gap width versus time delay, the delay error is not shown on this graph since it is too small to be seen, ±3 ps, and the spatial error is a result of the energy fluctuation of the laser system although the weak scaling between energy and blast wave velocity, $E^{1/4}$, means that the energy fluctuations had minimal effect on the blast wave trajectory. Fig. 5.5 shows the velocity at which the cluster disassembly wave is seen to be travelling, also shown
on this figure are the kinetic energies of ions and electrons travelling at these speeds.

While it is not possible to retrieve the electron density in the regions where the two heating beams cross, with single beam interferometry, looking at the electron density at the right hand edge of the vertical blast wave far from the intersection we can see the extent of the radiative precursor in front of the blast wave. The measured electron density of the vertical shock is shown in Fig. 5.3 and it can clearly be seen that the observable radiative precursor of the blast wave only extends \( \sim 0.5 \) mm ahead of the shock. So, given that the gap between the two blast waves is so much larger than the observed precursor extent (1 mm versus 0.5 mm) ahead of the shock that it cannot be the responsible mechanism for cluster disassemble. Even if there were some longer range radiative precursor, which does not cause sufficient ionisation of the upstream medium to be directly observed using our interferometer, this precursor would move at the same velocity as the shock. The observed velocity of the disassembly wave is far in excess of that of the shock. Given the poor match between the shock velocity and the disassembly wave velocity it seems unlikely that some previously unobserved secondary radiative precursor could be responsible.

![Figure 5.3: Measured electron density profile of the vertical blast wave far from the region where the vertical and horizontal blast waves cross. The shock is clearly visible with a precursor extending \( \sim 0.5 \) mm ahead of it. This density profile was obtained with \( \sim 500 \) mJ of laser energy over a duration of 1 ps in argon clusters from a modified Parker 99 valve with a backing pressure of 30 bar giving a neutral density of \( \sim 2 \times 10^{19} \) cm\(^{-3}\).](image)

If the blast wave itself is not responsible for the cluster-free gap then the only
other source for the cluster disassembly wave must be some form of emission from the initial laser driven ionisation of the clusters or the hot plasma filament this initially creates at early time. The emission resulting in cluster breakup must then take the form of either electrons, ions or photons.

Figure 5.4: Graph showing the separation between the centre of the vertical blast wave and the front edge of the horizontal blast wave on the left hand side, $d_L$, as a function of time delay. The error in spatial position is a result of shot to shot energy fluctuations from the laser system. The error in the time delay is so small that it cannot usefully be seen on the graph, it is $\pm 3$ ps.

To better understand the nature of the process driving the cluster disassembly wave we first consider a photon propagating through the ambient medium. While clusters exhibit strong absorption of laser radiation this only occurs above a threshold intensity of $10^{12}$ Wcm$^{-2}$, this is because the strong absorption results from transient field effects that couple to a dynamic resonance in the cluster rather than single photon events [36, 37, 38]. For this reason we need only consider the behaviour of single photons interacting with the clusters, such a photon would need to be of a high enough energy that it could have a significant impact on a single atom of argon. Given that the first ionisation level of an argon atom is at 15.8 eV the medium will be transparent below this energy. Therefore an optical or UV flash that was sufficiently bright to be able to disassemble the clusters would result in a wave propagating at near to the speed of light in a vacuum. For this reason we
consider photons with 20 eV or more of energy in the calculations which follow.

As photons above the ionisation energy propagate through the medium they will undergo a random walk, as they are sequentially absorbed and re-emitted. When considering the propagation speed of a disassembly wave driven by these non-ballistic photons, we model this random walk as:

\[ L = MFP \times \sqrt{N} \]  

(5.1)

Where \( L \) is the length through the medium the photons have propagated from their origin, MFP is the mean free path of the photons and \( N \) is the number of steps taken during the journey. We define the MFP of the photons as the length over which transmission falls below 50%. When determining the mean free path of the photons we consider two models: 1) mono-atomic gas absorption and 2) absorption in uniformly spaced atomic clusters.

Shown in Fig. 5.6 is the transmission fraction of photons over a range of energies through 400 \( \mu \)m of mono-atomic argon gas with conditions matching those in the experiment [106]. Using these MFPs we can determine the number of steps which must be taken at the respective energies to propagate to the edge of the gas jet. While emission above these test energies, of the order of 100 eV, has been observed from the laser irradiation of clusters, the MFP at these energies is so long that they
would not have been stopped within the gas volume and the disassembly wave would be propagating at close to the speed of light [41].

![Graph of transmission vs. photon energy](image)

Figure 5.6: Transmission of photons propagating through a uniform gas matching the Parker 99 valve’s conditions during this experiment [106]. Here the absorption is determined over a propagation length of 400 µm.

If we consider a cold system consistent with generating both high enough energy photons and not leaving a medium which is strongly enough ionised to be detected using Schlieren or interferometric imaging, then processes like Auger decay could be viable candidates. This would give recombination times of order 2 to 20 fs in argon under these conditions [107]. Using these recombination times we can estimate the speed at which the disassembly wave would travel through the medium. Shown in Table 5.1 are the speeds associated with the given photon energies’ MFPs. The reader should note that the longer 20 fs recombination time was used to determine the lowest possible speed in an attempt to match the observed data.

Under the conditions of this experiment we would expect to have clusters with a mean radius of $\sim 30$ nm at a density of $10^{22}$ cm$^{-3}$ within the argon cluster. This length and density is sufficient that the absorbed fraction of photons in the 10-50 eV range would be greater than 50% within a single cluster. Therefore when applying our uniformly spaced cluster model the MFP of the photons is approximately the spacing between the clusters in the medium. This model leads to a very short MFP
Table 5.1: Table showing the mean free path (MFP) and velocity of the predicted disassembly wave with respect to the speed of light in a vacuum for photons propagating through a mono-atomic argon gas [106].

<table>
<thead>
<tr>
<th>Photon Energy/eV</th>
<th>MFP/µm</th>
<th>Velocity/C</th>
</tr>
</thead>
<tbody>
<tr>
<td>20</td>
<td>60</td>
<td>35%</td>
</tr>
<tr>
<td>30</td>
<td>80</td>
<td>49%</td>
</tr>
<tr>
<td>40</td>
<td>450</td>
<td>97%</td>
</tr>
<tr>
<td>50</td>
<td>2250</td>
<td>~100%</td>
</tr>
</tbody>
</table>

of ~110 nm. Clearly this would require a great number of steps for the photons to propagate far through the medium. The recombination times in clusters have been found to be ~15 fs for a range of cluster sizes [108]. Using this recombination time we arrive at a velocity for the cluster-cluster photon disassembly wave of ~600 ms\(^{-1}\). Even if we assume that the MFP is strongly underestimated in this calculation, the MFP would need to be four orders of magnitude larger to match the disassembly wave speed. This suggests that cluster-cluster propagation of photons seems highly unlikely as the mechanism driving the disassembly wave.

The other obvious candidates for producing the “cold” cluster disassembly are electrons or ions produced by the initial laser driven ionisation of clusters. We can use the velocity of the disassembly wave to determine what the necessary energy of the electrons or ions would be to match the known disassembly wave velocity. Clearly from Fig. 5.4 the disassembly wave is slowing as it propagates through the medium and as such the energy of the electrons or ions responsible will also be dropping, as shown in Fig. 5.5. At early time the kinetic energy corresponding to the observed velocity for an electron would be 3 eV and for an argon ion ~200 keV. The initial energies of the electrons or ions can be compared to the predicted particle energies of the nano-plasma model in order to determine whether they are in agreement [37, 40, 48].

While the nano-plasma model has been shown not to fully predict the behaviour of laser-cluster interactions with pulse of < 1 ps [52], it does produce experimentally well confirmed predictions with ‘long’ pulse durations into high Z clusters. This matches well with this experiment where the 1 ps duration is ‘long’ and argon (Z=18) is of high enough Z. However when using very short pulses or low Z atoms, for example hydrogen, the model’s predictive powers fail. The nano-plasma model predicts that under the conditions in this experiment the peak ion energies would be ~100 keV and peak electron energies would be hundreds of keV.
The 3 eV electron energy associated with the velocity of the disassembly wave is far too low to match the predictions of the nano-plasma model. It has also been demonstrated in previous laser-cluster experiments that there are electrons free streaming through the medium with a velocity of $\sim 10^7 \text{ ms}^{-1}$ under similar conditions. This is two orders of magnitude faster than the early time velocity of the disassembly wave and one order of magnitude slower than the peak velocity of the electrons created by the laser-cluster interaction [55, 109]. However this simple analysis does not account for the interaction of the electrons with the clusters and the author is not aware of any work in the literature which investigates the interaction of clusters with electrons in this energy regime. For this reason the electrons cannot be absolutely discarded as drivers of the disassembly wave.

Given the reasonably good match between the ion energies predicted by the nano-plasma model and the observed velocity of the disassembly wave in the medium, within a factor of $\sim 2$, the ions would seem the most likely candidate as the primary driver of the disassembly wave. The gradual slowing of the rate of gap growth would also suggest a mechanism where there is a gradual loss of energy available to the mechanism driving the cluster disassembly. This could be attributed to multiple small angle collisions of the electrons or ions. The laser irradiated clusters are not of one size only and the broad cluster size distribution will produce electrons and ions with a range of energies.

The higher energy ions exist in such small numbers that they may not destroy a sufficiently large number of clusters, as they propagate through the medium, to affect the energy deposition of the second heating pulse. The lower energy ion tail will also propagate through the medium, however a significant number may never catch up to the peak energy ions since the lower energy tail will also be losing energy to the medium. Additional to the nano-plasma model’s predictions, the ion energies measured here also correspond well with other experimental measures of ion energies generated by atomic clusters [40]. A similar argument can also be applied to electrons propagating through the medium, this would further complicate separating the effect of the ions from the electrons.

Since the disassembly wave should affect the whole length of the vertical blast wave not just where the second heating beam is crossing the first, it should therefore be possible to use our interferometer to measure the free electron density in the region where the clusters have disassembled but the second heating beam has not interacted with the medium. Our interferometry set-up is capable of detecting an
Table 5.2: Table comparing predicted disassembly wave speed using different mechanisms at early time. Photon-Gas refers to 20 eV photons propagating through mono-atomic gas as compared to Photons-Clusters where the photons are being stopped by near solid density argon clusters. The ion and electron velocities do not take account of interactions with the clusters and are instead the vacuum velocities of the particles.

<table>
<thead>
<tr>
<th>Mechanism</th>
<th>Velocity/ms⁻¹</th>
<th>Percentage of C</th>
</tr>
</thead>
<tbody>
<tr>
<td>Photons-Gas</td>
<td>1.05×10⁸</td>
<td>35%</td>
</tr>
<tr>
<td>Photons-Clusters</td>
<td>600</td>
<td>0.0002%</td>
</tr>
<tr>
<td>Ions</td>
<td>6×10⁵</td>
<td>0.2%</td>
</tr>
<tr>
<td>Electrons</td>
<td>1.64×10⁸</td>
<td>55%</td>
</tr>
<tr>
<td>Disassembly Wave</td>
<td>1×10⁶</td>
<td>0.33%</td>
</tr>
</tbody>
</table>

average ionisation state as low as Ar⁺¹ at the atomic number densities used in this experiment. None of the interferometric data collected indicates ionisation of the upstream medium beyond the blast wave precursor. These findings are consistent with the work of Kagawa et al. on collisions between ions and clusters. Their work shows that argon ions with energies as low as 3 keV propagating through an argon cluster medium are capable of disassembling the clusters. Furthermore their work also shows that the medium then turns into a neutral monoatomic gas within 100 ps with little or no residual ionisation [110].

As is clear in Fig. 5.5 the disassembly wave, and hence the charged particles, are slowing as they propagate through the argon cluster medium. This is a result of electric and nuclear collisions between the particles and the ambient medium, which also lead to the disassembly of the clusters. The author can find no previous studies aimed at determining the stopping power of clusters for ions or electrons. A standard tool for modelling the ion stopping power of a wide range of materials is Ziegler’s Stopping Range In Matter (SRIM) code, which is based on Ziegler, Biersack and Littmark [111, 112, 113].

An attempt was made to create an intuitive model of a cluster medium to investigate the stopping of ions, here slabs of solid density argon, representing the clusters, were separated by regions of vacuum. To model a reasonable length of the medium many slabs were required. Unfortunately whilst the model on first inspection appears to be working well, it is in fact artefacts from the codes operation at the multiple boundary regions between solid and vacuum which are dominating the output. This is no surprise given the extreme number of regions through which an
Figure 5.7: Comparison of the experimental stopping power of the argon clusters and the predictions of SRIM for argon gas and argon ions with an initial energy of 220 keV. a) Graph showing the decrease in ion energy against time as it propagates through the medium. b) Graph showing the stopping power of the medium with distance, the reader should note that the gas density is dropping as we move away from the origin. Courtesy of H.W. Doyle

ion must be propagated in the intuitive slab-vacuum model compared to a few layers of thicker material which the code was intended to deal with. Intuitively we would still expect clusters to have a higher stopping power due to their high local density and the field enhancements which can occur in clusters [52].

Instead of trying to create a detailed model with the clusters included, we used a simple gas model. We matched the gas density in the simulation to the real density profile of the Modified Parker 99 valve starting 2 mm from the nozzle, see Sec. 4.3.2, and launched 220 keV ions through the medium. This gives a decreasing density gradient as the ions propagate outwards, modelling the region to the left of the vertical blast wave.

As can be seen in Fig. 5.7a the overall trend of the ion energy loss is well modelled by the SRIM monoatomic gas simulation at early time. However the clusters clearly have a greater stopping power than the monoatomic gas as can be seen in Fig. 5.7b. It is not surprising that the stopping power is consistent with neither monoatomic gases or solids since the cluster medium’s behaviour is somewhere between that of solids and gases. As such it is likely that the dynamics of any shielding effects occurring in the clustered gas are substantially different from those of solids or gases.
5.4 Additional Observations.

When the horizontal to vertical delay time becomes sufficiently large there are no clusters left in the region through which the horizontal beam then propagates. Despite the poor energy absorption of monoatomic gas ($<5\%$), the horizontal beam will deposit energy as it propagates through the un-clustered region. This creates a narrow pre-heated region in front of the vertical blast wave. If a blast wave were to propagate through a pre-heated region it would pick up additional energy as it collects mass into the thin shell of the blast wave and could then be accelerated by it.

To quantify this behaviour we would wish to determine the free electron density of the preheated material, however to do so would require vertical interference fringes. This would satisfy the conditions of the Abel inversion in terms of the image window needing to return to zero free electron density. The intensity of the horizontal heating beam as it passes through the un-clustered medium will still be $\sim 10^{17}$ W cm$^{-2}$ [114]. At such an intensity the literature suggest that the argon atoms will be ionised to the Ar$^{+7}$ state [115]. Using this ionisation state combined with a gas number density of $3 \times 10^{-18}$ cm$^{-2}$ we arrive at a free electron density $\sim 2 \times 10^{19}$ cm$^{-2}$. It is however clear that where the horizontal beam intersects the vertical blast wave there is an electron density enhancement even if it cannot be easily quantified. It has been suggested that this mechanism could be used to introduce a temperature perturbation at specific points along the shock. This perturbation could then potentially be used to seed the Vishniac overstability [10].

5.5 Summary.

The data produced by this experiment shows that from the initial ionisation of the cluster medium a ballistic wave is launched which results in the disassembly of the clusters and very little, if any, residual ionisation. From the measured electron density profiles and velocity consideration the shock and precursor can be discarded as the source of the ballistic wave. The mean free path (MFP) of photons of $\geq 20$ eV generated in the laser-cluster interaction are then considered since they possess enough energy to result in re-emission from the medium without causing significant ionisation which would have been detected by our interferometry diagnostic. Photons of 50 eV are shown to have a MFP so long, in a mono-atomic argon gas model,
Figure 5.8: Here the horizontal beam has been delayed by 7.9 ns and the second blast wave will have formed off the edge of the images. a) Dark-Field Schlieren image, here the distortion of the vertical blast wave is clearly visible where the horizontal heating beam crosses as a bump on the LHS and also on the RHS. b) Shows the phase map from the interferometry on the same shot. It is clear from the phase map that there is an electron density enhancement where the horizontal heating beam crosses the vertical shock. Abel inversion is not possible due to symmetry considerations.

that they almost immediately exit the system and therefore would drive a prompt disassembly wave propagating close to the speed of light. Lower energy photons moving in a random walk are shown to theoretically propagate two orders of magnitude faster than the measured ballistic wave. An intuitive cluster to cluster model is also introduced but again it indicates that photons are not the most likely drivers of the ballistic disassembly wave.

The nano-plasma model code is then deployed to predict the peak energies of ions and electrons generated by the laser-cluster interaction. While the predicted velocity of ions from the interaction match well with the ballistic disassembly wave and the electrons appear too fast, these predictions have not sought to account for the dynamics of the particle-cluster interaction. For this reason neither electrons or ions can be discarded as the driver of the ballistic wave. The propagation of ions is then investigated using the SRIM code and it was shown that the code does not model the stopping power of clusters for ions accurately. An attempt to use multiple slab-vacuum elements in the model highlights a weakness in the SRIM
code that cause its results to be dominated by un-physical boundary effects in this model.

Clearly it is very helpful in terms of building a laboratory astrophysics platform that the cluster medium does not persist in front of the blast wave and hence does not alter the dynamics of its evolution. However, further studies must be carried out to verify these results in other species and to determine whether electrons or ions are the primary drivers for cold cluster disassembly. In particular the outcome of a hydrogen cluster experiment where there is no radiative precursor at all should be pursued since with heavier atoms even with horizontal fringes the precursor will prevent effective diagnosis of the ionisation state close to the blast wave.

In principle the role of electrons can be separated from the ions by exposing the system to a sufficiently large magnetic field. When the field is applied the electrons and ions will start on a curved path and with a large enough field the electrons will begin to gyrate about the field lines only slowly drifting outwards. Using this mechanism it should be possible to effectively freeze the electrons such that only the ions, on an altered path, are still propagating through the medium and interacting with the clusters. Also in such a system photons would continue unaffected. This experiment could then be used to test all three possible mechanisms. However it does require large magnetic fields, which the author estimates to be \(\sim 300\). Which cannot at this time be generated without destroying the equipment used to generate it.

Furthermore this experiment demonstrates the need to carefully tailor the density profile of the cluster medium, using a skimmer for example, when attempting to use it as an imaging ion source. Studies of the stopping power of clusters would be best carried out using a particle accelerator for ions or an electron gun to control the bandwidth and peak energies.
Chapter 6

Astra-Gemini Experiment: Blast Wave Results.

In this chapter the experimental campaign using the Central Laser Facility’s Astra-Gemini Laser during August and September of 2012 is described. Particular focus is given to the experimental results pertaining to the radiative blast wave data in this chapter. Results gathered on the laser-cluster interaction obtained on this campaign are discussed in Chap. 7. The experiment was carried out by the author and a team comprising members from Imperial College, the University of Oxford, the Atomic Weapons Establishment and the Central Laser Facility. Analysis of the optical streak data and Schlieren imaging was conducted by the author and the analysis of the interferograms was conducted by the author and D. Bigourd. This work represents the first cluster based experiment ever conducted on the Central Laser Facility’s Astra-Gemini Laser.
6.1 Experimental Setup.

The experimental campaign made use of both beam lines of the Astra-Gemini laser, here on referred to as Gemini. One beam line was used to drive the experiment and the second beam line to derive a second harmonic optical probe beam. Two gas jets were used to generate the atomic clusters, the AASC jet and the Peter Paul valve with large bore nozzle which are described in Sec. 4.3.2. However unless otherwise stated the Peter Paul valve has been used due to operational issues with the AASC jet. The heating beam is focused, using an f/20 parabola, into the centre of the cluster volume. The driving beam and the second harmonic optical probe pass through the interaction region perpendicular to each other. Similarly the streaked Schlieren back-lighter also passes through the interaction region perpendicular to the heating beam. Shown in Fig. 6.1.

The single time frame probes were at 400 nm and the streaked Schlieren back-lighter was at 532 nm. The Be filters block all photons with <1 keV of energy setting the threshold for all images captured using the x-ray pinhole cameras. Both pinhole cameras used on the Astra-Gemini campaign of Andor DX420s as detectors. The single frame imaging used Andor iXon cameras for the late time images and Allied Vision Technologies Stingray cameras for the early time images. The streak line uses a doubled Nd:YAG laser and a Hadland streak camera with high sensitivity CCD camera.

6.2 Optical Streak Camera Data.

6.2.1 Optical Streak Camera Set-Up and Technical Considerations.

The optical streaked Schlieren data was gathered on a Hadland Imacon 675 streak camera with a high sensitivity CCD camera imaging the phosphor screen, see Sec. 3.1. The sweep rate was set to 5 ns/mm on all images shown in this thesis. The back-lighter pulse was generated using a Continuum PowerLite q-switched doubled Nd:YAG laser together with a rail holding a series of flat glass plates which produces multiple small reflections. If temporal delay between the small reflections is properly set and spatial overlap this effectively stretches the pulse temporally by stacking many smaller pulses together.

The interaction region was imaged using four 50 cm anti-reflection coated lenses arranged as a 1 to 1 imaging system using a 1 mm Schlieren stop. A schematic of
Figure 6.1: Schematic of the optical streak camera set-up used on the Gemini campaign. Where PBCs are polarising beam cubes, WPs 1/2 waveplates, Ms mirrors and Ls 50 cm AR/AR lenses. The probe pulse was centred at 532 nm delivering \(~20\) mJ of energy over a duration of \(~60\) ns.

this arrangement is shown in Fig. 6.1, the probe system delivered \(~20\) mJ on target and had pulse duration of \(~60\) ns. The spatial extent of the probe pulse is only sufficient to image one half of the blast with a spatial extent of \(5\) mm. The spatial resolution of the streak camera system was determined at the centre of interaction region to be \(28.7 \ \mu m\). The optical self emission of the laser-cluster interaction will be very short lived, delivering what can be considered a delta function signal to the streak camera. Therefore the duration of that signal can be treated as giving the temporal resolution of the streak camera set-up, which gives a temporal resolution of \(~0.9\) ns. Shown in Fig. 6.2 is an example of the stretched YAG pulse used as the back-lighter for the streak camera.

As is clear in Fig. 6.2 the back-lighter has an uneven spatio-temporal profile which will naturally lead to an uneven streak image. However because the electron density is so much greater at the shock in a Schlieren image it is possible to normalise each individual time slice of the Streak Schlieren Image (SSI) using the intensity at the shock front as the divisor. Shown in Fig. 6.3 is an example of a streaked Schlieren image which has not been normalised (A) and SSI which has been normalised (B). This method does not fundamentally alter the data but it does however make the data much easier to visualise.

However on the higher pressure and energy shots the precursor generated by the shock has a large spatial extent and also leads to a significant free electron density
Figure 6.2: Streak camera image without the Schlieren stop in place showing the temporal and spatial profile of the optical back-lighter pulse. This pulse is \( \sim 60 \) ns long in duration and has a spatial extent of 5 mm. The reader should note the hard edge at the top of the image which shows the location of the upper edge of the streak camera slit.

Figure 6.3: Example of streaked Schlieren images (A) before and (B) after each time slice in the image has been normalised with respect to itself. This is an early time shock launched in 12.7 bars of krypton with 12.76 J. The fine scale structure in (B) between 10-12 ns is an artefact of the probe pulse not the shock.
ahead of the shock. Given the uneven intensity of the probe pulse this can lead to the precursor producing a higher signal level than the shock itself. This can clearly interfere with the tracking of the trajectory of the blast wave, in these cases it is necessary to manually “remove” the precursor during the analysis. This is easily achieved in the majority of images since the free electron density down stream of the blast wave is very low giving almost no signal, and the blast wave’s free electron density is extremely high compared to the near precursor such that the blast wave can be easily distinguishable from the precursor.

Additionally the image is smoothed using a moving point average over a $3 \times 3$ cell to smooth high frequency noise and make the image easier to visualise. This smoothing of the images has an effect of $<0.5\%$ on the measured deceleration parameter, as discussed in Sec. 6.2.3, taken from the images and as such is an acceptable modification to make. The position of the shock is found by fitting a Gaussian distribution to each temporal slice of the streak and noting the position of the peak. The error in this position detection is taken to be the half max half width of the fitted Gaussian.

Taking the log-log of the trajectories retrieved in this way gives the deceleration parameter of the blast wave. The theory underpinning this is described in detail in Sec. 2.3.3. The reader should note that a non-radiative blast wave will have a deceleration parameter of 0.5 and the shock will be slowing due to its expansion. Any value less than 0.5 indicates that the shock is slowing faster than under its’ own geometry and this indicates that the shock is radiative. Any value above 0.5 indicates that the shock is accelerating with respect to the non-radiative case and energy is being fed into the system.

The early time plasma filament must evolve into a blast wave shell structure before the data can be analysed from a Sedov-Taylor model. Shown in Fig. 6.4 is a 6 ns window with the laser arriving at $\sim1/2$ ns. Here the blob of early time plasma can be seen before it evolves into a blast wave. This shows that data gathered in the first $\sim4$ ns after the initial interaction should not be treated as a blast wave and it is only after this point that analysis from a Sedov-Taylor like model can reliable begin [5, 104].

### 6.2.2 Early Time Velocity and Precursor Extent.

At early time the velocity of the blast waves in this experiment is high, however it is dependent on both drive energy and the number density of the local medium.
Figure 6.4: Streak Schlieren image in a 6 ns window with the laser pulse arriving at \( \sim 0.5 \) ns. This image shows that a Sedov-Taylor like blast wave does not develop until an interval of \( \sim 4 \) ns has passed from the laser-cluster interaction. In the first 3 ns of this image we can see a thick plasma shell formed from the laser-cluster interaction which can then be seen to form the thin shell structure which is a fully formed blast wave. This image was normalised and captured in krypton at a backing pressure of 12.7 bar and a laser energy of 12.76 J.

Shown in Figures 6.5 and 6.6 are the mean velocities of blast waves in the first 20 ns after their formation, which is 4 ns after the laser-cluster interaction itself, in krypton and xenon respectively.

As expected increased drive energy results in a higher shock velocity while increased gas jet backing pressure reduces the shock velocities. This is consistent with the increased density of the downstream material which the shock is propagating into. This behaviour is observed in both krypton, Fig. 6.5, and in xenon, Fig. 6.6. Furthermore the shock velocities are lower in xenon than in krypton for a given backing pressure. This will result both from the higher mass density of the xenon at a given pressure, and also the higher temperature of the xenon precursor slowing the propagation of the shock. The error shown here is estimated from the spatial and temporal resolutions of the streak camera set-up.

The streaked Schlieren images also provide a diagnostic of the blast wave pre-
Figure 6.5: Graph showing measured the relationship between shock velocity and drive energy over a range of pressures in krypton. As anticipated the shock velocity is higher with increased drive energy, however the higher backing pressure shoots also have a higher velocity which is contrary to the expected behaviour.

Figure 6.6: Graph showing the relationship between shock velocity and drive energy over a range of pressures in xenon. As anticipated the shock velocity is higher with increased drive energy and lower velocities on are measured for higher density shots.

cursor extent. The reader should note that while the signal incident on the optical streak camera is a convolution of the free electron density of the shock and precursor structure with the probe beam. Due to the inhomogeneity of the probe pulse we cannot use this diagnostic to measure the form of the precursor however it does still given an approximate measure of the spatial extent of the precursor.

This measurement is limited to fairly early in the blast wave evolution since at
Figure 6.7: Shown here are four time frames in krypton (left) and four time frames in xenon (right). In both sets the images were captured at 5, 10, 15 and 20 ns respectively after the laser-cluster interaction. The data was captured in krypton and xenon at 37.3 and 22 bar respectively with a drive energy of ∼13.5 J. These backing pressures match the mass densities in the krypton and xenon as closely as possible. Here, as expected, the radiative pre-cursor is longer in the xenon and in both cases is very long.

later times the precursor is obscured by the upper edge of the slit on the streak camera. Shown in Fig. 6.7 are four time steps captured in krypton and xenon at 37.3 and 22 bar respectively with a drive energy of ∼13.5 J. The time steps shown in Fig. 6.7 are at 5, 10, 15 and 20 ns after the initial laser-cluster interaction. These time steps are integrated over 1 ns in accordance with the temporal resolution of the streaked Schlieren set-up. Here the radial propagation of the blast wave is projected as propagation from right to left in the line-out. The backing pressures were selected to match the mass densities in the krypton and xenon mediums as closely possible.

The left hand edge of the line-outs in Fig. 6.7 is the position of the upper edge of the slit. Since the laser-cluster interactions are spatial centred on the streak camera slit we can therefore measure the maximum radius to which a blast wave can expand before falling outside the diagnostic window, this a maximum measurable radius of 2.81±0.03 mm.

The backing pressures used in Fig. 6.7 produce strongly radiative shocks and
Figure 6.8: Shown here are four time frames in krypton (left) and four time frames in xenon (right). In both sets the images were captured at 5, 10, 15 and 20 ns respectively after the laser-cluster interaction. The data was captured in krypton and xenon at 24.3 and 16.5 bar respectively with a drive energy of $\sim 13.9$ J. These backing pressures match the mass densities in the krypton and xenon as closely as possible. Here, as expected, the radiative pre-cursor is longer in the xenon and in both cases is very long.

spatially extended precursors. Therefore it is no surprise here that the precursor has reached the edge of the diagnostic window after only a matter of 15 ns for the krypton and 10 ns for the xenon respectively. However if we drop the backing pressure of the gas jet the radiative precursor will be reduced as the shock itself becomes less radiative. This is borne out in the data as shown in Fig. 6.8 where the same time steps for a krypton and xenon shock have been launched in 24.3 and 16.5 bar respectively. Here again the backing pressures have been chosen to keep the mass density of the ambient medium as close as possible in both systems.

We can confidently say that the spatial extent of the precursor shrinks as the backing pressure of the gas jet is reduced despite the inhomogeneity of the probe beam. We can also say that while the probe beam will distort the appearance of the precursor such that its shape is unknown it is still reasonably stable shot to shot. Therefore it is also valid to say that a change in signal level at any given point along the precursor does indeed represent a change in the actual free electron
Figure 6.9: Shown here are four time frames in krypton (left) and four time frames in xenon (right). In both sets the images were captured at 5, 10, 15 and 20 ns respectively after the laser-cluster interaction. The data was captured in krypton and xenon at 24.3 and 16.6 bar respectively with a drive energy of $\sim 4.57$ J. These backing pressures are match the mass densities in the krypton and xenon. Here, as expected, the radiative pre-cursor is longer in the xenon and however unlike the higher energy shots the pre-cursors here are more modest although still substantial.

density gradient at that point. So comparing the data in Figures 6.7 and 6.8 we are not only observing a shortening of the precursor extent with reduced backing pressure but also a fall in the absolute free electron density gradients at all points along the precursor.

However as is clear in Figures 6.7 and 6.8 the precursor grows so fast that it falls outside the edge of the streak camera slit very early in the evolution of the shock. As is evident when dropping the energy, as shown in Fig. 6.9 with krypton and xenon at 24.3 and 16.6 bar respectively driven with $\sim 4.57$ J, the precursor shrinks enormously. This indicates that the precursor spatial extent can only be properly tracked in the lower energy shots. Presented in Figures 6.10 and 6.11 are the measured spatial extents over a range of pressures in krypton and xenon respectively at $\sim 4.5$ J measured from the shock to the end of the observed precursor.

The low pressure data in krypton has not been included since it is unreliable. The signal produced by the shock and precursor is so low that it is difficult to separate
Figure 6.10: Graph of the spatial extent of the shock precursor at four time steps in krypton over a range of backing pressures with a drive energy of $\sim 4.5$ J. The pre-cursor extent can be seen to grow in the first 15 ns and then drop off after 20 ns.

Figure 6.11: Graph of the spatial extent of the shock precursor at four time steps in xenon over a range of backing pressures with a drive energy of $\sim 4.5$ J. The pre-cursor extent can be seen to grow in the first 15 ns and then drop off after 20 ns. The reader should note that the final measure of precursor extent at 16.6 bar is not included because it extends beyond the streak camera slit.

The reader should note that the final measure of precursor extent at 16.6 bar is not included because it extends beyond the streak camera slit. However it is clear that in both the krypton and xenon the precursor extent reaches a peak length before falling off again. This is consistent with energy loss in the system as a whole and hence a reduction of emission from
the shock into the precursor. The size of the precursor’s measured here are very large. Given that, for example, the shock shown at 10 ns in Fig. 6.9A, which was captured in krypton at $\sim 4.57$ J, the FWHM of the shock in the Schlieren image is $210 \pm 28 \mu m$. Comparing this FWHM the precursor is an order of magnitude larger. A further discussion on the shock width and shock precursor can be found in Sec. 6.3.

6.2.3 Evolution of Blast Wave Trajectories.

The primary purpose of using the streaked Schlieren set-up was to capture data on the temporal evolution of a shock on a single shot basis. From this data it is possible to extract the deceleration parameter of a shock, as discussed in Sec. 2.3.3. Shown in Fig. 6.12 is a streaked Schlieren image with the position of the shock superimposed onto the image in black. All images analysed here are only tracked for 35 ns because a single streak is only back lit for approximately this period of time. Furthermore the higher the density of the ambient medium the shock is propagating through the stronger the observed signal is, thereby making it easier to track. It is possible to combine multiple streaks despite the energy variation shot to shot, $R \propto E^{1/4}$, however at this time only the earlier time streaks have been investigated.

The streak shown in Fig. 6.12 runs from the moment of the laser-cluster interaction, 0 ns, and is 25 ns in length. The position of the shock is determined by fitting a Gaussian curve to each individual time slice and selecting the peak as the position of the shock. This method is preferred over selecting the brightest point in the time slice since it prevents high signal noise spikes from distorting the tracking. The error bar shown in Fig. 6.12 are given as the half max half width of the fitted Gaussian [5]. In the region between 13 and 15 ns the probe pulse is particularly poor which results in erroneous signal with very large error bars. This feature is repeated in many of the streaked Schlieren images from this run and when extracting trajectories such regions are discarded to prevent distortion of the deceleration parameter calculation. In future this issue can be overcome by using a smoother brighter back lighter source, for example a modern optical pump laser.

As described in section 2.3.3 the deceleration of the shock should be governed by a power law and the coefficient of this power law is termed the deceleration parameter, $n$. The value of $n$ can most easily be determined by taking the log-log of the radius-time data and then fitting a linear trend in this space, where the gradient of this linear fit corresponds to $d(ln(R))/d(ln(t)) = n$. The gradient of this slope is
Figure 6.12: Streaked Schlieren image of an early time blast wave with the extracted shock position, with error bars, superimposed in black onto the original image. The region between 10 and 15 ns is an artefact of the probe pulse and should be disregarded. This shock was launched in krypton at a backing pressure of 10 bar with a drive energy 13.11 J.

measured over a 2 ns period in the figures presented in this thesis. However since we are tracking behaviour in the temporal domain the duration of the measurement window is important. Shown in Fig. 6.13 are the changing $n$ values of a shock using 3 different time bases, of 1.5, 2.0 and 2.5 ns respectively, for the fitting. As can be seen in Fig. 6.13 changing the time base does not alter the temporal position or measured deceleration parameter significantly. It is clear, as expected, that the shortest time base gives sharper features and the longest time base smooths the deceleration parameter. It is also important to note that the position and sign of the peaks and troughs in the data are insensitive to these changes in the time base. This highlights that they are not artefacts of the data analysis technique we have chosen. The time base of 2 ns was selected since it is longer than the temporal resolution of the streak camera system, $\sim$0.9 ns, and it is shorter than the dynamics we wish to observe, $>3$ ns [13, 14].

It should be noted that there are 8 camera pixels within the temporal resolution,
∼0.9 ns. For this reason the observed signal is a convolution of the streak camera function and the signal. Therefore it is not subject to Nyquist-Shannon theorem limitations since it can be considered that the signal is not being converted from a continuous to a discrete function [116]. When the deceleration parameter is calculated the signal is digitised, however the rate at which this occurs is up to the user. In this thesis the calculation has been performed at every camera pixel, such that a measurement has been performed every 125 ps of the data set. However the data is blurred at every camera pixel by the finite temporal resolution of the system. A sample error from the images analysed here can be as high as ±60%. However the reader should note that the algorithm for determining the error will be refined further and with the improved tracking of the shock front this error will likely be reduced.

The streaked Schlieren image from which Fig. 6.13 was extracted is shown in Fig. 6.12 where the shock is driven with 13.11 J in krypton at a backing pressure of 10 bar. The reader will note that in Fig. 6.13 there is no data shown between 12 and 19 ns which corresponds to the region in the streak image which is heavily distorted by the inhomogeneity of the optical back lighter probe pulse shown in Fig. 6.12. The error bars shown here are given by the \( R^2 \) value of a simple linear regression to the log position and log time data [117].

In previous experiments using laser-cluster interactions to launch radiative shocks an apparent velocity domain oscillation was observed when driving a shock in krypton above 11 J of incident laser energy with a number density of \( \sim 10^{19} \text{ cm}^{-3} \). The oscillation period in these experiments was found to be between 7-9 ns [5, 13]. This measurement has been shown to be consistent with the theoretical work of Chevalier and Imamura [12] by Skidmore [14]. Detection of this phenomenon is achieved by viewing oscillation of the n value, first the deceleration parameter will drop below \( n<0.5 \), a regime where the shock is losing energy, and then rises to \( n>0.5 \), a regime where the shock is gaining energy. The reader should note that the period of any observed oscillations must be near constant over a number of cycles in order to fit with Chevalier and Imamura’s theoretical work. It should also be noted that, as the shock expands, the geometry dictates that there is always a reduction in velocity, so rather it is the rate at which the shock slows which changes.

The physical mechanism which underpins Chevalier and Imamura’s work [12] is based on the behaviour of the shock cooling region. Where the “cooling” region is the region immediately behind the shock front, an example of this is shown in
Figure 6.13: Graphs showing the changes in the deceleration parameter over the first 33 ns of the blast wave’s evolution. This data is extracted from a shock in krypton at 10 bar with 13.11 J of drive laser energy. The streaked Schlieren image these values are extracted from is shown in Fig. 6.12. The time bases over which \( n \) is measured is 2.5, 2.0 and 1.5 ns respectively in these graphs.

Fig. 6.24. As the cooling region is heated by the hot material which the shock propagates into, the velocity of the shock is increased by the additional pressure behind it. However in parallel with this the rate of cooling will also increase, leading to a drop in the cooling region temperature and hence a reduction in the shock velocity. The shock will then begin this process again by incorporating more pre-heated upstream material and heating the cooling region again. A key feature of this instability is the invariance of the periodicity of the oscillation.

It is not clear in Fig. 6.13 that there is a real oscillation in the deceleration parameter however looking at Fig. 6.14 where the backing pressure, again in krypton and at 10.24 J, has been increased from 10 bar to 37.3 bar, an oscillation of the deceleration parameter does indeed seem to appear. The period of this oscillation is \( \sim 3 \) ns under these conditions. Furthermore, shown in Fig. 6.15 is a late time streak taken in Krypton with a backing pressure of 37.3 bar and drive energy of 10.66 J; this measurement continues on from the end of that shown in Fig. 6.14. Oscillations seem to appear again on this streak but with a slightly reduced oscillation period of \( \sim 2.5 \) ns. These oscillation periods are very close to the measurement period
and at this stage are treated as an indication that an oscillation is present rather than a definitive confirmation. Shown in Figures 6.16 and 6.17 are the streaks from which the apparent oscillations are extracted, as is clear in these figures there are regions with significant error in the blast wave tracking. These regions must be carefully removed from the calculation of the deceleration parameter adding further uncertainty to the measurement.

Figure 6.14: Graph showing the changes in the deceleration parameter between 13 ns and 33 ns of the blast wave’s evolution. This data is extracted from a shock in krypton at 37.3 bar driven with 10.24 J of laser energy. The streaked Schlieren image these values are extracted from is shown in Fig. 6.16.

The apparent oscillations, with stable periods, observed in the ∼10 J high backing pressure krypton shots is not repeated in the xenon data. Shown in Fig. 6.18 are the n values in the first 33 ns of evolution of a xenon shock, driven with 13.61 J of laser energy at a backing pressure of 22 bar. The mass density of 22 bars of xenon roughly corresponds to ∼34.5 bar of krypton and Fig. 6.14 is also the highest available energy in krypton. The shocks from which Fig. 6.14 and Fig. 6.18 are derived, are approximately comparable. In Fig. 6.18 there may be a single full oscillation between 12 and 15 ns. However, it does not show sufficient oscillations, with a stable period, that we would believe the velocity domain oscillation to be present. This is as expected since the level of radiation generated by xenon shocks
Figure 6.15: Graph showing the changes in the deceleration parameter between 34 ns and 50 ns of the blast wave’s evolution. This data is extracted from a shock in krypton at 37.3 bar driven with 10.66 J of laser energy. The reader should note that this a late time measurement. The streaked Schlieren image these values are extracted from is shown in Fig. 6.17.

Figure 6.16: Streaked Schlieren image captured in 37.3 bar of krypton with a drive energy of 10.24 J. It is from this image that the evolution of the deceleration parameter shown in Fig. 6.14 is extracted. The reader should note that the shock precursor has been manually removed from this image during the analysis in order to track the position of the shock more accurately.

is believed to be too high and hence drowns out any velocity domain oscillation of this type [5, 118]. The large errors and many features that can be seen in this data
Figure 6.17: Streaked Schlieren image captured in 37.3 bar of krypton with a drive energy of 10.66 J. It is from this image that the evolution of the deceleration parameter shown in Fig. 6.15 is extracted. The reader should note that the shock precursor has been manually removed from this image during the analysis in order to track the position of the shock more accurately.

Figure 6.18: Graph showing the changes in the deceleration parameter over the first 35 ns of the blast wave’s evolution. This data is extracted from a shock in xenon at 22 bar driven with 13.61 J of laser energy. The streaked Schlieren image these values are extracted from is shown in Fig. 6.19.

make it difficult to be certain. However the trajectory certainly is not a smooth and simple Sedov-Taylor blast wave.
Figure 6.19: Streaked Schlieren image captured in xenon at a backing pressure of 22 bars with a drive energy of 13.61 J. It is from this image that the evolution of the deceleration parameter shown in Fig. 6.18 is extracted. The reader should note that the shock precursor has been manually removed from this image during the analysis in order to track the position of the shock more accurately.

Shown in Fig. 6.20 is the evolution of the deceleration parameter in the 38 bar region in krypton at the lowest drive energy, 4.27 J. Unlike the higher drive energy counterpart in the 38 bar region there is no clear oscillation of the deceleration parameter in this time range. However it is possible that oscillations have not yet begun at this stage of the blast wave’s evolution, or that they are lost in the poor data region between 10 and 16 ns.

Previous experimental studies have also shown that there exists a cut-off energy below which the velocity domain oscillation will not be seen. This cut-off falls somewhere between 7 and 11 Joules for cluster blast waves of this type [5]. The data at ∼10.2 J in krypton with 37.3 bar of backing pressure also appears to show an oscillation after the first ∼20 ns, as can be seen in Fig. 6.14. This would suggest that the threshold for the onset of the velocity domain oscillation may fall below 10.2 J. However this data does suffer from the inhomogeneity of the probe beam, which prevents a proper analysis of the earlier time data where the onset of a velocity domain oscillation may be present. Furthermore the large errors in the data make it difficult to identify where there is definitely no oscillation.
Figure 6.20: Graph showing the changes in the deceleration parameter over the first 27 ns of the blast wave’s evolution. This data is extracted from a shock in krypton at 38.2 bar driven with 4.27 J of laser energy.

Figure 6.21: Streaked Schlieren image captured in 38.2 bar of krypton with a drive energy of 4.27 J. It is from this image that the evolution of the deceleration parameter shown in Fig. 6.20 is extracted. The reader should note that the shock precursor has been manually removed from this image during the analysis in order to track the position of the shock more accurately.

As previously mentioned the precursor must be removed from the streaked Schlieren images to prevent the inhomogeneous probe pulse from interfering with tracking of the shock front. It is possible that this technique can be refined further, improving the quality of this tracking and hence reduce the errors in deceleration
parameter calculation. Further refinement of the fitting algorithm may enhance the accuracy of the calculated values shown in Figures 6.13 through 6.20.

However this data indicates that there may be an oscillation with stable period present in the deceleration parameter. This behaviour is closely linked to the species, backing pressure and drive energy of the observed shock. It also indicates that the threshold drive energy for velocity domain oscillation may fall lower than previous measurements taken using similar gas jets but with a longer, ∼1 ps, drive pulse [13]. It is possible that the threshold for oscillation may fall somewhere between 7 and 10 Joules depending on the details of the laser energy absorption phase.

6.3 Second Harmonic Optical Probe Data.

The plasma was probed using both interferometry and Schlieren imaging to capture the electron density profile of the shock and the position at a given time. An attempt was made to split the probe beam by polarisation in order to capture two time frames on each shot. However due to the relatively low level of the polarisation purity of the available optics and the demands of a high-bandwidth sub 50 fs probe pulse this method was abandoned and a single time frame was captured of the plasma. The probe pulse is generated from the second beamline of the Astra-Gemini laser and before doubling to 400 nm will have ∼100 mJ of energy such that each of the four probe pulses would contain at least 20 mJ of energy, which is more than sufficient for any imaging required.

The plasma was viewed using a 1 to 1 imaging system using four 50 cm AR/AR coated achromatic lenses. A Mach-Zehnder interferometer was utilised for the electron density profile measurement and the Schlieren imaging is derived from a beam split from the interferometry line, Fig. 6.22 shows the optical arrangement of a single arm of the imaging system.

6.3.1 Free Electron Density of the Blast Waves Launched on Astra-Gemini.

The probe pulses have a peak intensity >10^{10} which given the large amount of material in the system as the beam enters and exits the chamber results in significant levels of B-integral [91, 119]. This non-linearity introduces phase front distortion to the laser pulse which interferes with the interferometry. The large spatial extent of the radiative precursor, as discussed in Sec. 6.2.1, which on some interferograms
extends to the very edge of the image can also prevent proper Abel inversion, c.f. Sec 3.1.2. When these two factors are combined it increases the difficulty of retrieving the free electron density profiles. For these reasons the analysis of these images is still on going and an example is included here for the sake of completeness. The probe beam includes a variable dogleg to control temporal overlap and the final turning mirror of the reference arm was automated to optimise spatial overlap.

Shown in Fig. 6.23 is an example of a raw interferogram captured in krypton at 37.3 bar of backing pressure in the Peter Paul valve with a laser drive energy of 10.24 J, 30 ns after the laser cluster interaction had occurred. In this image extreme distortions of the phase front can be seen in the region marked with a white circle, although the reader should note that there are other distorted regions in this image, this distortion is so severe that free electron densities cannot be extracted from it.

Careful analysis of such an image can still retrieve the free electron density of the shock. Shown in Fig. 6.24 is the free electron density profile measured from the interferogram shown in Fig. 6.23, with krypton at 37.3 bar driven with 10.24 J. This density could be extracted because care was taken to avoid the distorted regions.

As expected the free electron density profile shown in Fig. 6.24 shows a steeply
Figure 6.23: A raw interferogram of the blast wave launched in krypton at 37.3 bar of backing pressure in the Peter Paul valve with 10.24 J, 30 ns after the laser-cluster interaction has occurred. Within the white circle is a region modified by the phase front distortion from non-linear effects during beam propagation.

Figure 6.24: Retrieved free electron density from the interferogram shown in Fig. 6.23. This shock was launched in krypton at 37.3 bar with 10.24 J and was captured 30 ns after the laser-cluster interaction. The reader should note the large extent of the precursor, $\sim 600 \mu m$ ahead of the shock front.

rising shock front and a strong precursor with both long spatial extent and significant electron density. The precursor here is shorter than the precursors measured using the streaked Schlieren diagnostic Sec. 6.2.2. However the interferogram was generated 30 ns after the laser-cluster interaction, given that the precursor extent
peaks ~15 ns after the laser cluster interaction and then starts to reduce, the extent of the precursor measured here is consistent with these early time measurements. The interferometric measurement is also more precise than the streaked Schlieren images and accurately captures the shape of the shock and precursor.

It is clear from this analysis that over a range of conditions we will be able to extract the free electron density profiles of the shocks. This will allow us to set boundary conditions in the on going trajectory simulations being carried out using the HELIOS code by R. Scott of the CLF.

6.3.2 Schlieren Imaging of Shocks Launched on Astra-Gemini.

Schlieren imaging was also deployed on the Astra-Gemini run not because it was expected to deliver data of special interest but rather because it is a relatively simple diagnostic to implement and it can also be used to benchmark the shock position in order to confirm the reliability of other diagnostics. However rather than observing only the thin shell structure of the shock front a large amount of self emission was also observed along the laser axis in the vicinity of 400 nm. The emission is observable because it is not produced by a collimated source but rather comes from a number of point sources which are not blocked by the Schlieren stop. The self emission is so strong that the shock typically falls below the dynamic range of the camera. This feature is, we believe, a product of the laser-cluster interaction and not the shock driven through the ambient medium. For this reason it will not be discussed further here and is instead presented in Sec. 7.1.1, as is more appropriate.

6.4 Summary.

This chapter discussed the data relating to the dynamics of the shocks’ launched using a laser-cluster interaction gathered using the Astra-Gemini laser during the August-September 2012 campaign. A variety of diagnostics were deployed in order to capture the dynamics of the shocks launched from the hot plasma created by the laser-cluster interaction. We start with a measure of the early time velocity of the fully formed shock, such that comparisons can be made with other shock experiments, based on the trajectories tracked by the streaked Schlieren diagnostic. Then the spatial extent of the shocks’ radiative precursor is measured, again using the streaked Schlieren technique. This demonstrates the temporal evolution of the shock precursor, in particular it can be seen that the xenon shocks radiative precursor
peaks \(\sim 15\) ns after the laser cluster interaction. A similar although less distinct pattern is repeated in the krypton data. Ultimately the limiting factor on these measurements is given by the height of the imaging slit, this indicates that the precursor observed during this experimental campaign are unprecedented in their spatial extent.

Then the temporal evolution of the deceleration parameter of the shock is measured. While the error bars on this first iteration of analysis are quite large, it seems that there may be an oscillation with stable period in the deceleration parameter observed on some krypton shots at higher densities. In the xenon data no oscillations with stable period can be clearly identified. The observed velocity domain oscillations are believed to be driven by changes in the cooling region temperature behind the shock front as described by Chevalier and Imamura [12]. As previously observed on cluster blast wave experiments using the Vulcan laser there appears to be a drive energy threshold below which the oscillation is not observed [5, 13]. That threshold value was previously determined to be somewhere between 7 and 11 J, this range may in fact be narrower, falling somewhere between 7 and 10 J of drive energy. However the large errors on this measurement and the temporal duration of the oscillations as compared to the measurement window introduce the possibility that we are observing an artefact rather than a real oscillation, further work must be performed to improve the extraction method to tackle these issues.

We then move on to discuss the single time framed imaging of the shock structure. As of the writing of this thesis analysis for the interferometric retrieval of the free electron density of the shock is still on going, hampered by phase front modulation due to non-linear effects in the probe pulse and the spatial extent of the shock’s radiative precursor. However one example of the shocks free electron density profile is shown. From this measurement it is clear that these shocks do not deviate significantly from those generated on similar experiments in the past. Finally the single time frame Schlieren imaging is discussed, while this technique failed to produce a benchmark of shock position it did generate other interesting results which are discussed further in Sec. 7.1.1.

The reader should note that while several of the diagnostics fielded on this experiment did not yield the quality of data obtained on previous experiments with the CLF’s Vulcan laser [13] or at Imperial College [104] this was the first cluster blast wave experiment carried out on the Astra-Gemini laser system. Use of an ultra-short 400 nm probe pulse proved to be problematic and there is also clearly room for improvement to the streaked Schlieren back-lighter source.
Chapter 7

Astra-Gemini Experiment: Laser Cluster Interaction Results.

In this chapter we discuss in detail some of the results produced during the laser-cluster campaign the author participated during August/September of 2012 on the Central Laser Facility’s Astra-Gemini laser system. This chapter focuses on the data which relates to the laser-cluster interaction rather than the blast wave (c.f. Chapter 6) which evolves from this interaction. The experiment was carried out by a team comprising members from Imperial College, the Central Laser Facility, Oxford University and AWE. The author was involved with the experimental operation of all diagnostics. Both the AASC jet and the Peter Paul valve with large bore nozzle, which are described in Sec. 4.3.2, were used in the experiments in this chapter. However unless otherwise stated the Peter Paul valve has been used. The author conducted the analysis of the post interaction beam diagnostics, the x-ray pinhole camera data and streaked Schlieren data. The analysis of the x-ray temperature measurement using Ross filters was carried out by H.F. Lowe.
7.1 Post Interaction Laser Diagnostics.

Downstream of the interaction region either a ground glass plate or a pick off mirror to an optical spectrometer were present during the experiment. These diagnostics where used to quantify the changes that the remaining laser pulse had undergone during propagation through the cluster medium. These two diagnostics were only available with the Peter Paul valve while shooting Krypton & Xenon.

7.1.1 Absorption Measurement.

Approximately 1.5 m downstream of the interaction region a ground glass plate was placed in the heating beam path. This was imaged onto an Allied Vision Technologies Marlin CCD camera and attenuated with a set of neutral density filters. The energy fraction absorbed by the clusters can be determined by comparing the signal spatially integrated across the image between a cluster shot and no gas shot. A set of neutral density filters are used to prevent saturation of the camera. To the authors’ knowledge no previous work has measured significant backward or side scatter from the laser pulse as it propagates through the cluster medium and as such an energy measurement downstream of the interaction region should accurately capture the energy absorbed while the laser pulse propagates through the cluster medium [55].

Figure 7.1: Images of the ground glass with 12 J of energy into Krypton backed with 10 Bars of pressure. (A) shows the image when the spectrometer pick off mirror is out of the beam, (B) shows the pick of mirror steering light to the spectrometer and (C) shows a no gas shot. The reader should note that these images are not saturated and have been enhanced to show the beam structure.

Initially the signal is integrated across the ground glass plate images for the no gas cases. By fitting this data it was confirmed that there is a strong lin-
ear relationship between the observed signal and the measured input beam energy. Furthermore the equation of the linear fit make it possible to determine what the signal would be for a no gas shot even when no data corresponding to this exact value is available. The equation of the fitted line is given by \( \text{Energy} = (0.1354 \pm 0.0027) \text{Signal} + (-23.897 \pm 0.478) \), which is used to convert the measured signal with gas into an energy measurement. The ratio of the measured energy upon the laser drive energy then is the energy fraction transmitted through the clusters.

Fig. 7.1 shows three images of the ground glass. In Fig. 7.1A the pick off mirror has been moved out of the beam, in Fig. 7.1B the mirror is in place to acquire the transmitted spectrum and Fig. 7.1C shows a no gas shot. Clearly when the mirror is in the beam path it will reduce the measured signal by blocking a fixed fraction of the pixels. If the structure of the transmitted beam is not excessively complex then it is possible to determine the factor by which the mirror alters the signal level. By using this correction the size of the data set was increased. This extra data improves the quality of the surface plot and highlights the strong laser energy absorption properties of clusters.

Figures 7.2 and 7.3 shows the absorbed fraction of the laser energy as 3D surface plots of energy versus backing pressure for Krypton and Xenon respectively with the medium being generated by the Peter Paul valve. For the Krypton the lowest absorption is still \( \sim 65\% \) and the highest is at \( \sim 97\% \) corresponding respectively to high energy low pressure and low energy high pressure regimes. The Xenon is a slightly stronger absorber with a minimum absorption of \( \sim 75\% \) and a maximum of \( \sim 97\% \) again corresponding to high energy low pressure and low energy high pressure.

As is clear in the Xenon surface plot, Fig. 7.3, in the pressure region between 20-25 Bar the absorption reaches a maximum and forms a plateau even at high input energies, \( \sim 14.5 \text{ J} \). However when the pressure is increased further the absorption starts to drop off, particularly at the higher energy end of the spectrum. The behaviour of the Krypton clusters, Fig. 7.2, is similar although the drop in absorption for high pressure at the lower energies does not appear to be very pronounced.

In Section 2.2.4 the interaction between a laser field and an atomic cluster is briefly explained. However in brief, clusters become strong absorbers of laser energy when the incident field exceeds an intensity threshold, this threshold has been measured to occur at an intensity of \( 10^{12} \text{ Wcm}^{-2} \) where the field starts to resonantly heat the cluster [36, 37, 38]. Given the large size of the clusters in this experiment the
expansion of the clusters in the laser field will be dominated by hydrodynamic pressure within the cluster where energy is transferred through inverse Bremsstrahlung from field to cluster.

After the experimental run the Gemini pulse was found to have a significant pre-pulse $\sim 63$ ps ahead of the main pulse with an intensity of $\sim 10^{14}$ Wcm$^{-2}$. The energy in the pre-pulse is significantly smaller than that of the main pulse as can be seen in Fig. 7.4, which shows a log plot of the temporal profile of the Gemini pulse. Given that the energy absorption fraction shown in the measured data is so high that the pre-pulse cannot destroy or saturate the clusters prior to the arrival of the main pulse, this introduces a number of possibilities: A) a thin filament is bored through along the central axis of the beam but the majority of clusters in the focal volume are largely unaffected and so can still absorb a large fraction of the main pulse, B) the preheating of the clusters occurs without destruction or saturation of the clusters and C) the pre-pulse is strongly absorbed in the wings of the gas jet.

While the peak intensity of the pre-pulse is $\sim 10^{14}$ Wcm$^{-2}$ the wings of the focus
Figure 7.3: Figure showing 3D surface plot of the absorption by the atomic clusters of laser energy with respect to the drive energy and the backing pressure for Xenon in the clusters produced by the Peter Paul Valve.

will have a far lower intensity which, unlike the main pulse, will be below the $10^{12}$ Wcm$^{-2}$ threshold at which clusters start to strongly absorb laser energy. This would mean that the only region strongly affected by the pre-pulse would be close to the laser axis. If the pre-pulse deposited enough energy into the clusters within this filament it could lead to the destruction of the clusters thereby reducing the laser energy absorption of the main pulse in the laser filament. Such an interaction would be observable on the ground glass plate as a significantly brighter central spot that would persist shot to shot over at least a small range of energies and pressures. Such a behaviour is not observed and it seems unlikely that an intense pre-pulse but with relatively little energy would be able to saturate or destroy clusters of the size generated by the Peter Paul valve.

In the second scenario the cluster are pre-heated in a thin filament along the laser axis, similarly to the first scenario, however the clusters are not destroyed here or do not expand so far before the arrival of the main pulse that they cease to strongly absorb the laser energy. Even though the intensity of the pre-pulse is high enough to
enter the strong absorption regime of the clusters it still carries only a small amount of energy. This means that the large clusters produced by the Peter Paul valve will not have a large hot population of electrons which will limit the expansion rate of the cluster. The reader should note that the cluster expansion time is not strongly dependent on the initial cluster radius [43].

In all scenarios the clusters in the wings of the jet will absorb some energy from the laser pulse. However in the third scenario the pre-pulse is absorbed entirely by the clusters in the wings of the gas jet, with the main pulse then being absorbed by the clusters in the body of the jet. While the third scenario is the least problematic to understand, and would allow the main pulse to deliver the maximum energy to the clusters closer to the centre of the jet, it seems that this is not the case. If the pre-pulse penetrates beyond the edge of the gas jet a number of electrons will be liberated at the centre of the jet, these electrons could then undergo a number of physical processes which result in second harmonic emission. On the Schlieren imaging lines self emission was observed close to the centre of the jet as shown in Fig. 7.5, for Krypton over a range of backing pressures irradiated with 13 Joules of energy.

A limited set of transmission filters were used in the imaging system to try to
Figure 7.5: Self emission at 400 nm observed on the Schlieren imaging lines without the backlighter. All images were captured at \( \sim 13 \) J in Kr over a range of pressures: A) 38.2, B) 24.3, C) 10 and D) 6.8 Bar respectively. The clusters were produced with the Peter Paul valve and the laser propagates from the right.

determine the wavelength of the self emission filaments, the transmission curves are shown in Fig. 7.6. Only the IF400 filter, a narrow band interference filter centred at 400 nm, gave any signal on the cameras which suggests that the majority, if not all, of the self emission is at the second harmonic (400 nm) of the driving laser pulse.

When the backing pressure of the jet is increased the number density of the plume and the size of the clusters grows. Intuitively this would suggest that at higher backing pressure the pre-pulse should be consumed by the wings of the gas jet and as the pressure drops the pre-pulse deposits more energy in the centre of the jet. As shown in the images of Fig. 7.5 reducing the backing pressure increases the signal level observed on the camera which is consistent with the intuitive description of the pre-pulse absorption in the cluster jet. Here two of the three possible scenarios are mixed and dependent on the backing pressure of the jet.

This intuitive description does give a coherent reason as to why electrons would
be liberated from the cluster prior to the arrival of the main pulse, but crucially due to the large size of the clusters and the small amount of energy in the pre-pulse the cluster temperature is unlikely to be so high as to disassemble the clusters before the main pulse arrives.

However given the complex nature of the cluster medium it is difficult to determine the interaction which leads to this emission and it is likely that there are a number of contributing mechanisms. These could range from the interaction of the free electrons with the positively charged clusters and free ions all the way through to electrons passing through the surface of near solid density clusters [120, 121].

Previously a model has been produced to explain 3rd harmonic emission from an atomic cluster where a resonance exist between the driving field and the electrons of the clusters, which exist in a dual population of hot electrons and a cold electron core, [122]. However this model was designed to deal with a different class of experiment where the clusters are much smaller than the wavelength of the driving laser field, which starts to break down in the large clusters produced by the AASC jet and the Peter Paul valve. Further the optics used in this experiment were made of BK7 which has extremely poor transmission at 267 nm, the 3rd harmonic of the driving field [123]. For this reason it is possible that there was also emission at the 3rd that could not be observed due to the optical elements in the system.

However while the mechanisms at play are unclear some interesting features can be observed. Fig. 7.7 shows self emission captured as before in argon at a backing pressure of ~18 bar driven with an energy of ~13 J. The clusters in this image were
Figure 7.7: Self emission at 400 nm observed on the Schlieren imaging lines with the shock shell faintly visible above the self emission filament back lit $\sim$30 ns after the heating pulse arrives. This image was captured in argon at a backing pressure of $\sim$18 Bar driven with an energy of $\sim$13 J. The clusters were produced with the AASC jet and the laser propagates from the right. The colour scale runs from black through white onto violet, this scale is arbitrary and intended only to highlight the features of the image.

produced with the AASC jet which makes direct comparison with the Fig. 7.5 more difficult, however in both instances the clusters are large and in same size regime. Since argon does not cluster as readily as krypton the $\sim$18 bar argon shot has a mass density equivalent which falls between the 6.8 and 10 bar krypton shots, see Sec. 4.3.

In Fig. 7.5 the region of strongest emission is close to the laser axis as would be expected, since this will be the region with the highest free electron density, however we can see that there are cusp like structures extending perpendicular to the filament. Furthermore the cusp like structures appear to be directional, pointing toward the origin of the laser pulse on the lower side of the filament and toward the direction of laser propagation on the upper side of the filament. This feature is repeated in the xenon and krypton data captured with the Peter Paul valve. The origin of the
cusp like structure and the directionality are both unknown. An azimuthal field generated by electrons propagating along the laser axis would not result in this form of directionality in electron motion. The reader should also note that this image is not captured in a plane parallel to the floor but rather the optical probe lines form an X shape through the interaction region. So the images are captured at $\pm 45^\circ$ to the plane parallel to the floor. This indicates the complexity of the process at work here, which remains unclear at this time.

### 7.1.2 Transmitted Spectrum.

Downstream of the interaction region a broadband pick off mirror to an optical spectrometer was present during the experiment. It could be driven in or out of the beam to aid in absorption measurements described in Sec. 7.1.1. The spectra were measured using an Ocean Optics HR2C0905 compact spectrometer with Xenon & Krypton cluster gas targets. Shots were taken without gas to capture the spectrum of the pre-interaction laser pulse, an example of such a spectrum can be seen in Fig. 7.8.

![Spectral profile of the Astra-Gemini laser after the interaction region with no target gas present. The laser is centred at 800 nm with a spectral bandwidth of $\sim 35$ nm. The reader should note the intensity scale is in arbitrary units.](image)

When the laser pulse interacts with the clusters the majority of the energy is absorbed, see Section 7.1.1, however the measured transmitted spectra show significant modulation and broadening with a pronounced blue shift. It is known that a laser pulse propagating through an ionising medium will be blue shifted [124, 125]. Given that the clusters have been pre-ionised by the pre-pulse at 63 ps, see Fig. 7.4,
and the pedestal of the main pulse, which will be significant at least 10 ps before the peak of the pulse, it seems probable not only that the laser pulse was propagating through a plasma but also that there could be a significant ionisation gradient in the plasma. This gradient would be both spatial and temporal both of which can aid in the blueshifting of the transmitted laser pulse [126]. Since the nature of the laser-cluster interaction is unclear, particularly due to the pre-pulse, it is difficult to determine the exact mechanism driving the spectral features. Therefore a selection of spectra are shown, here over a range of conditions, in order to inform the reader of the broad trends observed, but without seeking to identify or explain the physical mechanisms which may underpin this behaviour.

![Figure 7.9: The post-interaction spectrum of the laser pulse. Here the laser had passed through Xenon clusters backed at 9.8 Bar with a range of energies. The respective drive energies in the three spectra were A) 4.06 J, B) 9.39 J and C) 13.11 J. Both spectra A and B are saturated due to inadequate filtering, however features can still be observed in these images.](image)

Fig. 7.9 shows the transmitted spectrum through 9.8 Bar of Xenon over a range of energies. As is immediately clear the spectrum has broadened significantly toward the blue for all energies ranging from ~825 nm to ~600 nm as compared to the pre-interaction spectrum which covers wavelengths from ~825 nm to ~775 nm, see Fig. 7.8. As can be clearly seen the higher energy spectra are saturated which results from the higher level of transmission through the lower backing pressure cluster.
Figure 7.10: The post-interaction spectrum of the laser pulse. Here the laser had passed through Xenon clusters backed at 16.6 Bar with a range of energies. The respective drive energies in the three spectra were A) 4.51 J, B) 10.22 J and C) 13.59 J. The spectrum is both blue shifted and modulated.

Figure 7.11: The post-interaction spectrum of the laser pulse. Here the laser had passed through Xenon clusters backed at ~29.2 Bar with a range of energies. The respective drive energies in the three spectra were A) 4.89 J, B) 9.91 J and C) 13.82 J. A peak can be seen to grow at ~675 nm as the laser input energy is increased.

medium, see Section 7.1.1 for details.

Fig. 7.10 shows the transmitted spectrum through 16.6 Bar of Xenon again over a range of energies. The change in backing pressure changes the spectral modulation
of the transmitted pulse, in particular the lower energy driving pulse produces a broader spectrum. However the overall form of the modulation does appear to be conserved at this backing pressure.

Figure 7.12: The post-interaction spectrum of the laser pulse. Here the laser had passed through Krypton clusters backed at 10 Bar with a range of energies. The respective drive energies in the three spectra where A) 4.03 J, B) 9.93 J and C) 13.67 J.

However at the highest available backing pressures, ~29.2 Bar, the transmitted spectra become significantly less complex, as can be seen in Fig. 7.11. However there is still broadening of the spectrum toward the blue with features extending down to ~670 nm similar to the lower backing pressure data. A feature can be observed growing from the background noise at ~675 nm as the drive energy is increased. It is also of particular note that in the spectral region covered by the pre-interaction pulse there are spikes at 785 nm and 825 nm and a plateau between 792 nm and 802 nm which is preserved across the full energy scan.

Looking at the Krypton spectra we can see in the lower backing pressure data (10 Bar), Fig. 7.12, that the spectra still show the characteristic blue shift. However the spectral region around 800 nm is depleted with significantly more signal present in the 650-700 nm than in the Xenon spectra. A more limited data set is available at higher pressures, 24.3 & 33 Bar respectively, in Krypton. These spectra, Figures 7.13 & 7.14, bear more resemblance to the Xenon spectra, in particular the 33 Bar spectrum is almost identical to the 29.2 Bar Xenon spectra.
Figure 7.13: The post-interaction spectrum of the laser pulse. Here the laser had passed through Krypton clusters backed at 24.3 Bar at a drive energy of 13.71 J.

This data indicates to some extent the complexity of the interaction between a laser pulse and cluster medium on both microscopic and macroscopic scales. Not only are the spectra heavily modified with strong blue shift but there are also strong variations with laser energy, atomic species and the backing pressure of the jet.

Figure 7.14: The post-interaction spectrum of the laser pulse. Here the laser had passed through Krypton clusters backed at 33 Bar at a drive energy of 12.91 J.
### 7.2 X-Ray Pinhole Images of Early Time Plasma.

An X-Ray Pinhole Camera (XRPC) was placed 3 cm from the interaction region perpendicular to the laser propagation direction. The XRPC had a pinhole diameter of 100 \( \mu \text{m} \) and was 30 cm long, giving a magnification of \( \times 10 \). The XRPC was filtered with 25 \( \mu \text{m} \) Beryllium giving a cut-off photon energy of \( \sim 1 \text{ KeV} \) [106]. A description of the function of an XRPC is given in Sec. 3.2. Only the early time plasma created by the laser-cluster interaction would be expected to reach temperatures high enough to emit x-rays which would pass through the Beryllium filter. This means that the radiative blast wave which evolves from the early time plasma filament will not be visible on the XRPC. Mounted on the back of the steel tube which forms the XRPC body was an Andor DX420 camera with a back illuminated uncoated CCD array and no coating, the spectral sensitivity of the Andor is shown in Fig. 7.15.

![Graph showing the spectral response of an Andor DX420 camera](image)

**Figure 7.15:** Graph showing the spectral response of an Andor DX420 camera used on the XRPC in order to image the early time plasma produced by the laser-cluster interaction. The DX420 used on this experiment used a back illuminated uncoated CCD array (BN). This graph was modified from [74].

As is clear in Fig. 7.15 the cameras’ response is good in the soft x-ray regime just above the Beryllium filters cut-off photon energy of \( \sim 1 \text{ KeV} \). Furthermore the reader should note that the temperature of the radiative shock launched by the laser-cluster interaction does not result in a high enough photon energy to pass through this filter. The pinhole camera is high magnification with a large pinhole,
therefore the resolution is dominated by the geometrical limit of the pinhole rather than diffraction or the detector resolution. The estimated resolution of the complete imaging system including the Andor camera was \( \sim 103 \mu m \).

Figure 7.16: X-ray pinhole camera backing pressure scan images taken in xenon with the drive energy broadly the same on all shots. With (A) 9.8 bar 13.11 J, (B) 16.6 bar 13.59 J and (C) 29.6 bar 13.82 J. The direction of laser propagation is from left to right in the image. The plasma filaments can be seen to narrow with increasing backing pressure and dip in signal level can also be seen along the laser axis.

Only the early time plasma will be high enough temperature to generate the photon energies required to penetrate the beryllium filter. This early time region is approximately 100 ps when the energy density can be estimated to be \( \sim 10^5 \text{ Jcm}^{-3} \) [5]. Therefore we would not expect the electrons to remain confined to the laser heating region but also to extend radially before the temperature falls to a level where the emitted radiation will not penetrate the beryllium filter.

Shown in Fig. 7.16 are a number of XRPC images captured in xenon. It is clear from these images that the backing pressure of the gas jet has a very strong impact on the width of the plasma filament and the brightness. As the backing pressure increases the filaments narrow from \( \sim 400 \mu m \) to \(< 300 \mu m \) for backing pressures of 16.6 bar and above with the maximum available drive energy. Furthermore, as can be seen in Fig. 7.16C, at the highest backing pressure of 29.6 bar the brightness of
the image increases rather than becoming narrower. At the centre of these filaments a lower signal region can also be observed, this will be discussed later.

Figure 7.17: X-ray pinhole camera backing pressure scan images taken in krypton with the drive energy broadly the same on all shots. With (A) 10 bar 13.67 J, (B) 24.3 bar 13.40 J, (C) 33 bar 13.54 J and (D) 38.2 bar 13.91 J. The direction of laser propagation is from left to right in the image. The plasma filaments can be seen to narrow with increasing backing pressure and dip in signal level can also be seen along the laser axis.

This pattern is repeated in the krypton XRPC images as shown in Fig. 7.17. Indeed the narrowing of the filament is even more pronounced between the 10 bar 24.3 bar krypton than the 9.8 bar and 29.6 bar xenon data, Figures 7.17 A and B, 7.16 A and C respectively. The 10 bar krypton images at 10 bar are so wide that they exceed the width of the pinhole camera’s CCD chip. Between 24.3 and 33 bar
the width of the filament falls from ∼400 µm to < 300 µm.

Figure 7.18: Raw interferograms and electron density profiles taken in 40 bar of argon at times of (a) and (d) 400 ps, (b) and (e) 11 ns and (c) and (f) 72 ns. Taken from [127].

It seems unlikely that this narrowing of the filament would be associated with self focusing in the clusters. Self focusing has been observed in both the long and the short pulse regimes [127, 128]. However to the authors knowledge self focusing has not been investigated using x-ray pinhole imaging but only using optical imaging techniques. These techniques show that over a range of pressures and drive energies the energy deposition in the clusters forms a number of pinches, as shown in Fig. 7.18 rather than a narrowing with backing pressure as observed in the XRPC images collected on the Astra-Gemini campaign.

The increased number density of the higher backing pressure shots would however limit the spreading of the filament in the first ∼100 ps due to collisions. This could easily account for the narrowing of the filament with backing pressure and also provides an insight as to why the filament does not narrow any further after 16.6 bar. The fact that there is no further narrowing with increased backing pressure also suggests that at this point the XRPC is giving an accurate representation of the effective focal volume within which the laser intensity is high enough to strongly
heat the clusters, c.f. Section 2.2.4. This analysis would indicate that the effective focal volume in both the xenon and krypton jets is \(< 300 \mu m\) when above 16.6 and 24.3 bar respectively. With the Peter Paul valve these backing pressures correspond to mass densities of \(1.81 \times 10^{-5}\) g cm\(^{-3}\) for the xenon and \(1.7 \times 10^{-5}\) g cm\(^{-3}\) for the krypton.

The reader will also note the filaments shown in Figures 7.16 and 7.17 change diameter along there length as well as with backing pressure. For example this is very clear in Fig. 7.17B where the filament expands from \(~200 \mu m\) to \(~300 \mu m\). The resolution of the XRPC limits the accuracy of this measurement, however it does show clearly a qualitative change in filament diameter with length.

When comparing the xenon filaments, Fig. 7.16, with the krypton filaments, Fig. 7.17, a clear difference emerges. In krypton the filaments expands from left to right, the direction of laser propagation, however in the xenon the filament thins from left to right. On first inspection this could again be attributed to self focusing of the laser pulse in the cluster medium, with the focusing occurring to the left of the image for krypton and to the right of the image for xenon. Not only does this answer seem too “convenient”, it also fits poorly with the data since the rate at which the diameter change occurs is not significantly affected by the backing pressure as is the case with self focusing in clusters [127].

A further feature which is observed in the XRPC images is a clear dip in signal at the centre of the filaments with a peak to peak width of \(~200 \mu m\). Fig. 7.19 shows the integrated signal over a 400 \(\mu m\) of the filament taken in Krypton at 13.40 J with a backing pressure of 24.3 bar along with the raw image. Understanding what processes are encoded in this feature can be highlighted by reference to limb darkening, where a uniform density sphere or cylinder of emitting material appears darker at the edges than it does at the centre [21].

We are observing emission along a chord of the object, since the x-rays will not be strongly absorbed by the ambient medium. At any given point on our detector the emission appears higher at the centre because there is more emitting material to observe there [129, 130]. However the inverse case where the edges appear brighter than the centre will occur when observing a shell like structure rather than a uniform bulk structure. Here when observing a chord through the shell we encounter more emitting material than when looking through the centre of the object.

This indicates that the dip observed at the centre of some of the XRPC images
Figure 7.19: XRPC image with an inset line out of integrated signal across the width of the emitting filament showing a dip in signal along the central axis. This shot was taken in krypton at a backing pressure of 24.3 bar and a drive energy of 13.40 J. The whole filament has a FWHM width of $\sim300 \, \mu m$ and the dip has a width of $\sim200 \, \mu m$ over a length of 400 $\mu m$ of the filament.

comes from the observation of a shell like structure rather than of a uniform bulk plasma. At first inspection such a feature could result from the formation of the shock. Shocks launched using atomic clusters have previously been shown to have a temperature of $\sim10$ eV, which would not produce emission that could penetrate the filter and generate the observed images [131]. Additionally for the filament created by the laser-cluster interaction not to be seen and then for the blast wave to been seen on the XRPC the initial temperature would have to fall below 1 keV and then rise above 1 keV due to the compression of material into the blast wave shell. This seems an improbable scenario given the apparent thickness of the shell structures observed on the XRPC.

This structure could be the result of electron blow out in the laser focal volume creating a shell structure rather than a bulk filament. If we consider a Gaussian focal spot then there exists a gradient in the electric field which points out of the
centre of the focal spot. This will exert a force normal to the propagation direction of the laser pulse on the free electrons. Since it is the behaviour of the electrons which is primarily responsible for emission from the material this would suggest that the XRPC images are showing a shell of electrons about the laser axis generated by the blow out of electrons from the centre of the focus of the laser pulse.

However this feature does not appear under all conditions, again as the backing pressure is increased the feature can be seen to shrink. Eventually the filament takes on the shape of a simple cylindrical structure as shown in Fig. 7.20 with 33 bar of krypton driven at 13.54 J. This pattern is repeated in the XRPC images captured with xenon clusters as is shown in Fig. 7.21 which shows the integrated signal over a 400 $\mu$m length of the filaments shown in Fig. 7.16 which have a drive energy of $\sim$13.5 J with backing pressures of 9.8, 16.6 and 29.6 bar respectively. In Fig. 7.21 the peak to peak width in (A) is $>$300 $\mu$m, in (B) it is $>$200 $\mu$m and in (C) there is no dip present.
This dependence on backing pressure can be explained by the increased number density around the laser heated volume at the higher backing pressure. At a threshold number density the medium has sufficient resistance to the motion of the electrons down the electric field gradient of the laser field that they cannot escape far from their origin. It is possible that there is still some electron blow out at the higher backing pressures with the narrower filaments, which simply produces a dip smaller than the detector resolution of 100 µm. However on the optical streak set-up used to diagnose the trajectories of the radiative blast waves, see Sections 3.1 and 6, which evolve from the early time plasma a feature has been observed which appears similar to the XRPC double peak images.

Figure 7.21: Integrated line out taken from the XRPC images shown in Fig. 7.16. All shots were taken in xenon at (A) 9.8 bar with 13.11 J, (B) 16.6 bar with 13.59 J and (C) 29.6 bar with 13.82 J. As is clear the filaments not only narrow with increased backing pressure but the central dip in signal is no longer present at the high backing pressure.

A large flash of visible emission can be seen on the streak camera at early time from the laser cluster interaction on all shots, irrespective of the back lighter. However the features observed with the XRPC are not replicated by this self emission but rather it is only visible when the interaction region is back lit. Therefore the data from the early time plasma captured on the streak camera is, as all Schlieren images are, a measure of the electron density gradients in that plasma. The spatial resolution of the streak camera system is 28 µm, which is substantially better than that of the XRPC camera which would allows us to discount the presence of a dip in the higher backing pressure shots.

Shown in Fig. 7.23 are two early time streak camera trace images where in (A) there is a back lighter and in (B) there is no back lighter both in krypton at 24.3 bar
Figure 7.22: Integrated line out taken from the XRPC images shown in Fig. 7.16. All shots were taken in krypton at (A) 24.3 bar 13.40 J, (B) 33 bar 13.54 J and (C) 38.2 bar 13.91 J. As is clear the filaments not only narrow with increased backing pressure but the central dip in signal is only present at the lowest backing pressure.

of backing pressure captured in 3.6 ns window. The early time shock can be seen as a bright streak in Fig. 7.23A with time going right to left. In Fig. 7.23 the signal level is approximately an order of magnitude lower in (B) with no back lighter than it is in (A) with the back lighter. In (B) a double peak feature can be seen in time, this is discussed further in Sec. 7.4.

By integrating across the temporal domain, around the self emission signal produced by the laser cluster interaction, we can observe the spatial form of the early time plasma. This region is the effective size of the temporal resolution limited window of this optical streak camera set-up which is ~1 ns. Fig. 7.24 shows the temporally integrated signal of the two early time streaked Schlieren images shown in Fig. 7.23. Clearly in Fig. 7.24A a pattern similar to that of the XRPC images, Fig. 7.19, can be seen. The peak to peak width of the feature at 24.3 bar is ~250 µm, which is reasonable consistent with the width ~200 µm found in the corresponding XRPC image.

When observing the higher pressure krypton streaks the double peak feature disappears following the same pattern as the XRPC data. This is shown in Fig. 7.25 for a shot at 13.68 J in 33 bar of backing pressure where a feature can be observed with a FWHM of ~1430 µm. This feature is very noisy which is likely to be a product of an increasing level of radiative self emission from the higher density krypton medium. As the reader will note the FWHM of the features captured on the streak camera is much greater than that of the XRPC. This is a result of two factors, firstly the data on the streak camera is gathered over a much longer time of
Figure 7.23: Streaked Schlieren images captured in a 3.6 ns window centred around the interaction time. In (A) the image is back lit and in (B) the image is not back lit. The early time shock can be seen in (A) as a bright streak with expanding radius and a double peak feature can be seen in (B). These streaks were captured in krypton at 24.3 bar with (A) 14.13 J and (B) 13.35 J with time going right to left.

∼1.2 ns rather than in the first ∼100 ps on XRPC and secondly the optical streak camera is not observing the actually plasma density but rather the free electron density gradients.

The data gathered on the optical streak camera is collected over a sufficiently long time for there to be significant expansion of the plasma filament. As stated in Sec. 6.2 the early time shock was measured to be travelling at ∼10 ms⁻¹, this would suggest that the early time plasma observed on the XRPC is travelling very fast indeed. The radiation generated by the krypton filament will also be heating the ambient medium. In fact the flash of optical emission shown in Fig. 7.24B indicates that the heating of the ambient medium is strong enough that the ambient medium produces significant emission of its’ own. This preheating of the ambient medium will liberate electrons around the laser heating region. Naturally the intensity of radiation will fall off as it propagates through the medium and so the number off
Figure 7.24: Temporally integrated signal of the first 1.2 ns after the laser cluster interaction. Both signals were captured in krypton at 24.3 bar, where (A) is a back lit image with 14.13 J drive energy and (B) is not back lit with a drive energy of 13.35 J. The reader should note that the signal in (A) is one order of magnitude larger than (B). The FWHM of the feature in (A) is $\sim 750 \ \mu m$ and the dip from peak to peak is $\sim 250 \ \mu m$ wide.

electrons will also fall off with distance from the laser heating region. This gradient would be visible on the optical streak camera but not visible on the XRPC since the preheated medium would still be, relatively, cold and not emitting at short enough wavelengths to penetrate the beryllium filter on the Andor camera.

The data gathered on the optical streak camera is in reasonably good agreement with the XRPC data. This strongly suggests that the laser-cluster interaction in the lower backing pressure regime results in the formation of a cylindrical shell of hot plasma well before the system evolves into a shock structure.

This work is in good agreement with earlier studies of non-local electron heat transport by Ditmire et al [109]. In this work it was shown with data gathered over a number of time frames after the laser-cluster interaction that a fast ionisation wave sweeps outward from the laser axis. This work includes XRPC images captured of argon plasmas which show filaments of radius 100 $\mu m$ where the cut-off energy of the filter is at 0.5 KeV. The region within which laser energy is deposited is $< 50 \ \mu m$ indicating that significant heat transport has occurred even at these very
Figure 7.25: Temporally integrated signal of the optical streak camera trace over the first 1.2 ns after the laser cluster interaction captured in krypton at 33 bar with 13.68 J. No dip in signal is visible at the centre of this feature and the FWHM is \( \sim 1430 \, \mu m \).

early times. However the quality of the focal spot used in those experiments was extremely high, allowing them to define carefully the region directly heated by the laser. This is not true of the Astra-Gemini laser where some 60\% of the laser energy falls outside of the \( 1/e^2 \) resulting in a much large and more poorly defined laser heating region. Additionally Ditmire’s experiments only used 300 mJ of laser energy deposited in a medium of relatively small atomic cluster whereas here 13 J of laser energy and very large atomic clusters were used. Even if this does not change the dominant mechanism underlying the observed data it is difficult to confirm that non-local electron heat transport has occurred here due to the size and quality of the Astra-Gemini focal spot.

### 7.3 X-Ray Temperature Measurements.

Three Ross pair filters were used to determine the signal level at three emission points to match to a Maxwellian temperature distribution as briefly described in Sec. 3.2 [76, 77]. Shown in Tab. 7.1 are the filter pairs used on the Gemini run along with their thickness. This filter pack was mounted on the front of an Andor DX420 camera with a back illuminated CCD array and no coating on the chip. A
snout was mounted in front of the filter and camera with a ∼1 cm opening at the end to prevent x-rays resulting from any accidental activation of material within the chamber from reaching the camera. Furthermore between the filter pack and the camera a 50 µm Be filter was mounted to safeguard against damage of the camera by scattered 800 nm laser light. The filter pairs Zn-Cu, Co-Fe and V-Ti are described as the “hard” (9.0-9.6 keV), “medium” (7.2-7.7 keV) and “soft” (5.0-5.4 keV) filters respectively, where the hard filter requires the highest photon energies to pass through the filters.

<table>
<thead>
<tr>
<th>Ross Pairs</th>
<th>Vanadium 8.1 µm</th>
<th>Zinc 14.7 µm + Iron 5.2 µm + Titanium 8.4 µm</th>
</tr>
</thead>
<tbody>
<tr>
<td></td>
<td>Titanium 12.2 µm</td>
<td>Copper 15 µm + Iron 5.2 µm + Titanium 8.4 µm</td>
</tr>
</tbody>
</table>

Table 7.1: Table showing the Ross pairs used to determine signal levels in narrow spectral bands allowing the identification of a Maxwellian temperature distribution of the laser-cluster plasma.

The reader should note that the intended thickness of the Co-Fe pair was 10 µm and 12.5 µm respectively rather than the 11.3 µm and 11.7 µm which the filter actually held. The filter thicknesses are close to their intended values and Tab. 7.1 shows the actually filter thicknesses as used. With a properly arranged Ross pair there is still some mismatch in the filtering leading to the so called “residual”. This residual represents an error on the measured signal level of any filter pair where harder and software x-rays than intended leak through. The thinner than intended Co-Fe pair lead to an increased residual level, which compromises the integrity of the measurement and therefore the matching to the Maxwellian temperature distribution. When designing the set of filter pairs it must be remembered that as the temperature of the plasma falls, any given set will fall further from the peak of the distribution and farther into the tail. This results in a flating of the signal which makes it difficult to fit the distribution well. The filter pairs must not be too distant from the peak of the distribution or the measurement will be poor quality. Here the filters have been selected and arranged to measure temperatures of >5 keV.

Using only the hardest and softest Ross pairs an estimate of the temperature can still be determined, albeit with a significant error as a first iteration. This data can then be re-processed using the estimated temperature to scale the signal from
the middle pair and then, repeating this iterative process. If this works the solution will converge toward the more representative value, of course still accompanied with significant errors. This methodology should not lead to large changes in the measured temperature from the hard and soft filters. It should, however, be able to confirm that the initial estimated temperature is valid.

Figure 7.26: Image captured through the Ross pair filter pack. Here, going right to left, the filters are ordered as Zn-Cu, Co-Fe and V-Ti, which are the hard, medium and soft filters respectively. This image was captured in Xenon at a backing pressure of 9.8 bar with 12.54 J of drive energy.

At the time of writing only the first iteration of this process has been performed, using the hard and soft filters. Included in this thesis are the estimated temperatures with error bound over a range of conditions. This data is included to give an impression of the probable temperature range of the plasma from which the shocks described in Chapter 6 were launched. Data was gathered in argon, krypton and xenon with both the Peter Paul valve and the AASC jet. However the data presented here will focus on the Peter Paul valve results because this matches with the conditions used to acquire the data described in Chapters 6 and 7. The error bars displayed on temperature data is simply the standard deviation over the data set and does not attempt to account specifically for experimental errors that are present in the measurement.

Shown in Fig. 7.26 is an image captured by the Andor camera where the outline of the filter pack can clearly be seen. The filter pairs, from right to left, are Zn-Cu, Co-Fe and V-Ti, running “hard” to “soft” with the image produced in Xenon clusters with a drive energy 13.54 J at a backing pressure of 9.8 bar.

Fig. 7.27 shows a graph plotting the effect of backing pressure on the measured temperature of the argon, krypton and xenon plasmas on high energy shots of \( \sim 13.6 \).
Figure 7.27: Graph showing the estimated temperature derived from the Ross pair data against the backing pressure in argon, krypton and xenon at \(\sim 13.6 \text{ J}\) of drive energy. The measured temperature can be seen to fall off with increased backing pressure. J. As is clear in this graph there is a reduction in temperature with increasing backing pressure. Since higher backing pressure corresponds to increased number density it could be hypothesised that a lower temperature is result of spreading the same or potentially slightly more total deposited energy over greater number of atoms. The reader should note that previous work suggests that the cluster should have initial temperatures of 5-10 keV immediately following the laser-cluster interaction [104].

The reader should exercise some caution when viewing this data, due to the selection of the filter pack the temperature measurement quality below \(\sim 5\) keV is not good and the quality will improve as temperature increases.

In Fig. 7.28 a graph of increasing drive energy over a range of backing pressures in argon is shown. In the 10 and 20 bar cases the temperature of the plasma increases as the drive energy increases, which can be attributed to more energy being imparted by the laser field to each cluster in the focal volume.

However Fig. 7.28 also indicates that on the highest backing pressure shots the temperature of the plasma only increases marginally. This might suggest that the clusters in the laser focal volume have been saturated with energy and therefore the temperature increase is damped. This should be accompanied by a decrease in the absorbed energy fraction measured behind the interaction region. Unfortunately
no data is available on the argon absorption with respect to energy and backing pressure, however shown in Sec. 7.1.1 is the data gathered in krypton and xenon. Both the krypton and xenon data indicate that above a threshold backing pressure (>30 and >25 bar respectively for Kr and Xe) the cluster absorption does indeed start to drop off, although it does remain high at >85%. This drop in absorption would explain a plateau of temperature in the argon clusters.

![Graph showing the estimated temperature derived from the Ross pair data against drive energy in argon over a range of pressures. This shows that the temperature increases as the drive energy is increased.](image)

7.4 Early Time Optical Streak Self Emission Data.

Observing the non-backlit optical streak Schlieren images at early time a bright flash of emission is observed. Since this feature can be observed when the streak is not backlit it must be an emission process resulting from the target medium rather than a free electron density gradient being measured by the normal Schlieren imaging system. Shown in Fig. 7.29 are four examples of this early time self emission, all these images are at full energy ~13.2 J with 16.5 bar and 22 bar of xenon (A and B) and 24.2 bar and 38.2 bar of krypton (C and D). These backing pressures were selected such that the mass densities in the xenon and the krypton are approximately
matched, they pair as A-C and B-D.

Figure 7.29: Early time optical streaks of the laser-cluster interaction self emission in xenon at 16.5 bar 13.59 J and 22 bar 13.14 J (A and B) and krypton at 24.2 bar 13.71 J and 38.2 bar 12.43 J (C and D). The pressures in xenon and krypton have similar mass densities in A-C and B-D.

The reader should note that the temporal resolution of the optical streak camera in this set-up is ~1 ns and so one must be careful when trying to observe and explain temporal features. The top and bottom of the self emission in the spatial direction
coincides with the position of the streak camera slit. Given the large spatial extent of the experimental volume around the interaction region and the strength of the signal observed here it seems reasonable to assume that the self emission region extends further than measured here.

The absolute signal levels between the xenon and krypton images should not be compared because they were taken on separate days and the streak camera alignment may not be consistent. However it is possible to compare between high and low pressure in a given species. On this basis it is clear that higher backing pressure shots generate more signal than the low and that the temporal behaviour of the signal is also affected by the backing pressure.

![Graphs showing the spatially integrated signal of the early time optical streak](image)

Figure 7.30: Graphs showing the spatially integrated signal of the early time optical streak, shown in Fig. 7.29, about the laser axis in (A) xenon at 16.5 bar 13.59 J and (B) krypton at 24.2 bar 13.71 J.

Shown in Fig. 7.30 is the spatially integrated signal for a 280 µm region around the laser axis over 4 ns in (A) xenon at 16.5 bar 13.59 J and (B) krypton at 24.2 bar 13.71 J. It is clear that under these conditions a double feature in time is created, this feature can also be seen by careful observation of the images shown in Fig. 7.29. Unlike the measurements taken with the long pulse back-lighter described in Sec. 6.2 the self emission nature of these measurement implies that the ratio of the peaks does carry some physical significance. The temporal resolution of the streak camera is ~1 ns and as such the two peaks fall just outside this resolution and can be resolved.

The distance about the laser axis was selected to match the spatial extent of
the earlier time signal, which gives a value of $280 \pm 14 \mu m$ for the radius of this feature, this is not wholly inconsistent with the hot filaments diameters measured in Sec. 7.2. An intuitive explanation for this feature is that the first flash of optical emission is from the laser-cluster interaction and is produced by the same mechanism as described in Sec. 7.1.1. The secondary flash, over the full length of the imaging slit, $\sim 1$ ns after the laser-cluster interaction then is product of some heating process. Given the speed at which this event must take place it cannot be attributed to the movement of charged particles that has been shown to disassemble the clusters, described in Chap. 5, and must rather be product of some other mechanism. The reader should note that on no gas shots, with all imaging beams in operation, there is neither a first or second peak observed on the streak. This would suggest that the observed signal is not simple scattered light from within the chamber but is instead dependent on the presence of a gas, however clusters are known to produce very little side scatter when irradiated by intense laser pulses further suggesting that this is an emission process rather than an artefact [37]. Furthermore the photo-cathode of this model of Hadland streak camera has an extremely low sensitivity at 800 nm, this also suggests that the observed signal is not scatter from the heating pulse. Given the low intensity of the probe pulses when compared to the level of filtering on the streak camera, composed of neutral density filters and a 532 nm interference filter and that the probe pulses arrived 30 ns after the laser-cluster interaction it also seems unlikely that the signal is the product of stray light from the short pulse probe lines.

The later time feature was not detected on the other optical diagnostics, where the early emission may share an origin with that shown in Fig. 7.5. This is most likely due to the higher sensitivity of the streak Schlieren line, the lower level of signal expected on this diagnostic from the outset means that it has lower ND levels than the single time frame diagnostics. Therefore the later time large spatial extent signal may fall below the detection threshold of the single time frame optical diagnostics which are typically deployed on such experiments [56].

### 7.5 Summary.

This chapter discussed the data relating to laser-cluster interaction experiments during the August-September 2012 Astra-Gemini campaign. Presented here was the first characterisation of the laser absorption fraction in a large atomic clusters with a short high energy laser pulse, using a ground glass transmission diagnostic. It is clear
that the majority of the laser energy is coupled into the clusters even in this extreme regime, which is consistent with previous measurements under different laser-cluster conditions [37]. This measurement has demonstrated a clear link between laser energy, atomic species and backing pressure of the cluster source. Furthermore it also provides an insight into the effect of the significant pre-pulse of the Astra-Gemini laser system, which was present during this campaign. Post interaction diagnostic studies of the drive laser spectrum have been presented which point clearly to the complex nature of the laser-cluster interaction with significant blue shift, possible due to ionisation in the medium, and strong dependence on the backing pressure of the cluster source.

Beryllium filtered x-ray pinhole camera images of the early time plasma filament generated in the laser-cluster interaction have been shown. These images show the variation not only in the width of the early time plasma with gas jet backing pressure but also the emergence of a cylindrical electron density distribution never seen before in a laser-cluster interaction experiment. This structure is further confirmed by investigation Schlieren streak camera images at early time. The early analysis of a Ross pair temperature diagnostic are discussed, indicating some broad trends in the early time temperature, which again depend on laser energy, atomic species and gas jet backing pressure. The early time plasmas were found to have temperatures of \( \sim 5 \) to 10 keV. Finally, the unexpected observation of the non-backlit early time streaked Schlieren images shows two flashes of self emission generated immediately after the laser-cluster interaction. We present a first interpretation of this result with reference to Zel’Dovich and Raizer analysis of an in air detonation of a nuclear weapon.
Chapter 8

Conclusions and Future Work.

In this chapter the conclusions of all the previous chapters will be drawn together followed by a discussion on how this work can and will be progressed in the future. This thesis has discussed laser-cluster interactions and the shock systems which evolve from the initial hot plasmas created by that interaction, as both a subject of merit in its own right and as a vehicle to enable high energy density laboratory astrophysics experiments.
8.1 Conclusions.

Chapter 1 is a short introduction to the motivation for the experiments contained in this thesis and seeks to set them in a broader context, with particular reference to laboratory astrophysics.

Chapter 2 sets out the theoretical underpinnings of this work. It explores the fundamental definition of a plasma and how we describe the behaviour of a plasma. Then we move on to laser-matter interactions, focusing on different target media, the behaviour of electrons in a strong laser field, the ionisation mechanisms at play in different intensity regimes and finally the interaction of a laser pulse with an atomic cluster, the target medium used in this thesis to efficiently create high energy density plasmas. Shocks are then broadly discussed from the basic piston driven shock to the Sedov-Taylor blast wave and evolution of radiative blast waves of the type investigated here. This description of shocks is then expanded to include the scaling methodology used to move between the laboratory and astrophysical scenarios via suitable choice of dimensionless parameters.

Chapter 3 moves on to a description of the measurement techniques used to diagnose different features of plasmas. This includes a description of how refractive index measurements combined with Abel inversion can be used to determine the free electron density profile of a shock (of suitable geometry) and its precursor at one time frame. Also Schlieren imaging is described as a method for detecting the position of the shock front. The optical streak camera is introduced which when combined with Schlieren imaging can be used to track the evolution of the shock trajectory with great accuracy on a single shot. Then a simple x-ray pinhole camera is introduced, which can be used to image hot plasmas early in time with good spatial magnification. Electron temperature measurements based on the x-ray emission spectrum are introduced in a regime where a Ross pair filter pack can be used to fit a Maxwellian temperature distribution.

Chapter 4 focuses on the key instruments and experimental methods used to create the high energy density plasmas described in this thesis. First the laser is described. This includes the simpler laser systems used extensively in their own right in the past and now as components in more sophisticated schemes, including q-switching and mode-locking. Moving on to the non-linear techniques used to generate the probe pulses used in the experiments and a brief description of the sources of damage in high intensity laser fields. Then chirped pulse amplification is introduced, which is of particular importance since all the laser systems used in the
experiments described here are built on this technique, with particular focus on the conditions for high fidelity re-compression of amplified laser pulses. After this the laser systems used in this thesis are briefly described along with the Glass beam line of the IC Cerberus laser system currently under construction and intended for use in future laser-cluster experiments.

Chapter 4 then moves on to the generation of the atomic cluster target medium. First the physical process that forms atomic clusters is described followed by a discussion of the semi-empirical Hagena scaling parameter before describing the three gas jets used in this thesis. The first of these is a modified nozzle Parker-Hannifin 99 valve used to generate the cluster target medium used in the cross beam experiment, Chap. 5 which has been in use on such experiments for a decade. Then measurements on the performance of the new AASC supersonic jet and the Peter Paul EH22G7DCCM with large bore custom nozzle are described. Both of these systems have proved to be valuable sources of large uniform volumes of large clusters at the \( \sim 1 \) cm scale.

Chapter 5 looks at data gathered on the Blackett Laboratory Laser Consortium Nd:Glass laser system using a two perpendicular beam geometry. Previously it was suggested that the clusters in the medium upstream of the shock were destroyed by radiation generated close to the shock front itself [104]. Discussions with the authors of this paper confirms that this was a only suggestion and is not supported by experimental or theoretical evidence. The work in this chapter demonstrates that the clusters are not disassembled by a emission from the blast wave itself, rather a ballistic wave is emitted from the early time laser-cluster interaction, which then disassembles the cold clusters without leaving any detectable residual ionisation of the medium. The velocity of this ballistic wave was measured.

Using intuitive models of the propagation of photons through the system it was shown that the ballistic wave is unlikely to be driven by the interaction of photons with the clusters. The only remaining candidates as drivers of the cluster disassembly wave are electrons or ions launched from the early time laser-cluster interaction. The nano-plasma model was used to estimate the peak energies of the electrons and ions launched under these conditions. While on first inspection the ions appear to match the velocity of the ballistic wave well, within a factor of two, and the electrons appear far too fast, this simple model does not account the interaction of the charged particles with the clusters in the ambient medium. Given the close fit of the ion velocity it is tempting to attribute the ballistic wave primarily to the ions, however the electrons cannot be discarded. A conceptually simple experiment is proposed
which should be able to determine whether the electrons or the ions are the primary
driver. Furthermore this experiment would also test conclusively whether photons
are the driver. The reader should note that this provides direct evidence of the
previously theorised position that there are no clusters immediately upstream of the
blast wave and so the radiation dynamics of the system should be that of a simple
mono-atomic gas. This adds support to the position that atomic cluster gases offer a
good environment for conducting scaled high energy density laboratory astrophysics
experiments.

Chapter 5 then briefly discusses the effect of the second heating beam at very
late time when the first shock has formed and no clusters are present to launch the
second shock. While the absorption is not strong enough to launch a shock it clear
that some energy is being deposited. This results in an acceleration of the shock
where the two heating beams intersect and a local enhancement of the free electron
density. It has been suggested that this mechanism could be used to create a tailored
medium intended to seed the Vishniac instability.

Chapter 6 presents data gathered on the 2012 Astra-Gemini campaign which
pertains to the shock launched by the laser-cluster interaction. This begins by de-
scribing the streaked Schlieren diagnostic and technical consideration when analysing
the data produced by this diagnostic. The trajectories recorded by the streak cam-
era are then used to determine the early time velocity of the shock, in order to aid
comparison to other shock physics experiments. The streaked Schlieren images are
then used to determine the extent of the radiative precursor in front of the shock.
This is of particular relevance since the large spatial extents of the precursors on the
high energy high density shots cause issues when trying to retrieve the free electron
density. It is shown that the precursor grows initially, peaking at $\sim 15$ ns after the
laser-cluster interaction before shortening again in both xenon and krypton.

Then the temporal evolution of the deceleration parameter is extracted from
the trajectories. As has been predicted theoretically [12] and shown previously
experimentally [13] under the correct conditions there will be an oscillation in the
deceleration parameter over time. This indicates a velocity domain oscillation of
the sort described by Chevalier and Imamura [12]. On high density krypton shots
there appear to be a few oscillations with a stable period of $\sim 3$ ns. However, in the
remaining data there is no clear null case without oscillations, but neither is there
a case with clearly periodic oscillations. Furthermore the region within which the
threshold for the onset of this instability to occur appears to fall between 7 and 10
J in krypton, which is lower than the previously observed 7 to 11 J. However due
to the high level of error in these measurements and long sample period it is not entirely clear that the oscillations are real, further development of the extraction algorithm may improve the quality of the measurements.

Chapter 6 then moves on to the 2nd harmonic probe data, starting with the experimental set-up which was used and issues originating from the modulation of the phase front from the accumulated non-linear phase. Due to this phase front modulation’s effect on the quality of the interferometric fringes used to retrieve the free electron density profiles and the radiative precursors long spatial extent interfering with the Abel inversion of the images this analysis is still on going. However an example of a free electron density profile retrieved from a lower density krypton shot has been presented which does show the well defined form of the shock and its clear radiative precursor. Finally single time frame Schlieren images are discussed, due to unexpected self emission from the early time plasma filament these images do not provide the accurate shock position data that was anticipated. However the self emission is of interest in its own right and is discussed in Chapter 7.

Chapter 7 presents all the data gathered on the 2012 Astra-Gemini campaign which pertains to the interaction of the laser pulse with the atomic clusters and features observed at early time before the shock forms. This chapter starts by investigating the characteristics of the post interaction transmitted laser pulse. This begins with a measurement of the absorbed energy fraction using a simple ground glass diagnostic. The data from this diagnostic repeats the previously observed high energy absorption fraction (>70%) of atomic clusters and confirms that short multi-Joule laser pulses can be used to drive high energy density physics experiments in atomic clusters of xenon or krypton. Furthermore a scan of drive energy and backing pressure has been performed to show the relationship between them and the absorbed energy fraction. Then the effect of the laser pre-pulse at 63 ps is discussed and the 2nd harmonic self emission of the laser heated clusters is explored with regard to the backing pressure of the gas jet and the laser drive energy. A pressure, energy and atomic species scan is shown of the transmitted spectrum of the laser pulse, these spectra include significant blueshift and with further modelling are expected to yield additional insight into the nature of the laser-cluster interaction.

Chapter 7 then moves on to the x-ray data gathered on the Astra-Gemini campaign. Using an x-ray pinhole camera and 25 µm Be filter it has been possible to image the hot early time plasma produced by the laser-cluster interaction. Here a relationship had been demonstrated between the backing pressure of the gas jet and
the width of the observed filament. Furthermore an unexpected dip in the signal at the centre of the filament is observed with its width again linked to the backing pressure. This feature is again observed in the early time backlit streak Schlieren images, the width of the feature is confirmed by both diagnostics. The origin of this feature is posited as non-local electron heat transport similar to that observed by Ditmire et al [109]. The chapter then moves on to consider the first iteration of an electron temperature measurement using a Ross pair filter pack. The temperature measurements follow the expected pattern and repeat previously measured results which indicate that the mean electron temperature of the plasma, created by the initial laser-cluster interaction, is $\sim$5 to 10 keV.

Finally in chapter 7 emission at 532 nm is measured on the optical streak camera. This includes an unexpected temporally resolved double peak feature in the first 2 ns after the arrival of the driving laser. The first flash in the initial 1 ns is likely linked to the 2\textsuperscript{nd} harmonic described earlier in this chapter on the single frame Schlieren imaging lines. It has a short spatial extent that seems consistent with the width of the filaments observed with the x-ray pinhole camera. The second emission event occurring after 2 ns has a large spatial extent spanning the full height of the slit on the streak camera. The speed at which this event occurs suggest that the medium is being heated by some form of photon emission.

8.2 Future Work.

The future evolution of this work will follow three distinct strands: 1) Laser development, 2) cluster source development and 3) experimental methodology development.

8.2.1 Laser Development.

Currently at IC the Cerberus laser system is under construction along with a dedicate cluster shock beam line, a schematic of this beam line is shown in Fig. 4.5. In its current incarnation the laser will deliver $\sim$2 J on target in a moderately short, $\sim$500 fs pulse, with an extremely high beam quality, similarly to the old BLLC Nd:Glass laser. Even in this incarnation it will allow for further development of the cluster launched shock experiments and the laser-cluster interaction experiments. However the laser system will soon be upgraded with an additional amplifier and larger compression gratings. This will slow the shot rate but increase the on target energy to 10 J or more. Having such a system in house will allow the IC group to
pursue new experiments in long campaigns of the sort not viable at laser facilities.

However the experiments described in Chapters 6 and 7 demonstrates that the laser-cluster interaction is an effective way of coupling energy into a medium which allows us to launch fast radiative shocks. In fact this technique is so efficient that it can rival experiments which use kilo Joules of drive energy. These other techniques have severe limitations on shot rate due to the thermal recovery time of the Nd:Glass based amplifiers. As of writing this thesis two companies, Thales and Amplitude, offer 30 J 30 fs laser systems with a reported shot rate of 1 Hz [132]. This thesis presents the first attempt at parameter scans with shock systems and laser-matter interaction measurements. This was made possible only by the relatively high shot rate of the Astra-Gemini laser a one shot every 60 seconds. However at an order of magnitude faster extremely fine scans of parameter spaces will potentially become almost trivial with these new laser systems.

8.2.2 Cluster Source Development.

Three distinct gas jets were used in the experiments in this thesis. In the past the Parker 99 was the staple in our work, however it is not ideal for some blast wave dynamics experiments due to the small volume of the gas plume and the steep density gradients within this volume. Additionally the small size of the viable target region means that the laser must be focused very close to the nozzle itself limiting the view angle available to the diagnostics and increasing the risk of damage to the jet. The application of the large bore Peter Paul valve and the AASC supersonic jet represent a significant change, in that the large and slowly varying volumes of large clusters which they produce provide a far better medium to conduct experiments in. However, at this time they cannot be cryogenically cooled, as the valve must be in order to generate clusters in hydrogen, which is of great importance since hydrogen is the non-radiative test case for many experimental scenarios.

At this time the Alameda Applied Sciences Corporation is developing a cryogenically cooled valve similar to that deployed during the Astra-Gemini campaign which will be able to operate at up to 10 Hz [133].

8.2.3 Experimental Methodology Development.

The experimental geometry used in Chapters 6 and 7 is extremely simple with a single beam and long focus through the cluster volume, Chapter 5 discusses an experiment with a more sophisticated geometry. In this experiment an enhancement
in electron density is observed on the early time shock. This introduces the possibility to controllable alter the shock front and seed spatial instabilities. This is of great importance since it could be used to explore the Vishniac instability [10], while avoiding all of the problems associated with the use of wire arrays or other physical blocks such as in Eden et al’s work [134]. Furthermore exploration of annular focusing optics also provide a chance to explorer colliding shocks such as that described in [6].

8.3 Closing Remarks.

This thesis demonstrates how the high laser energy absorption characteristics of atomic clusters makes them a unique and fascinating medium for high energy density physics experiments. The shocks in this experiment have exhibited, as previously observed only once before, the presence of a velocity domain oscillation of the shock front. Furthermore a wealth of data relating to the interaction of the high intensity pulse and large atomic clusters is presented. In the future by combining the high repetition rate of the new lasers systems that are coming into operation now with the high repetition rate large bore gas jets also coming to market as this thesis is written it will be possible to drive shock physics experiments with enough shots to make statistically significant statements about the observed data. This would represent a step change in how shock physics is conducted and the work presented in this thesis is a first effort in realising this new mode of operation.
References


