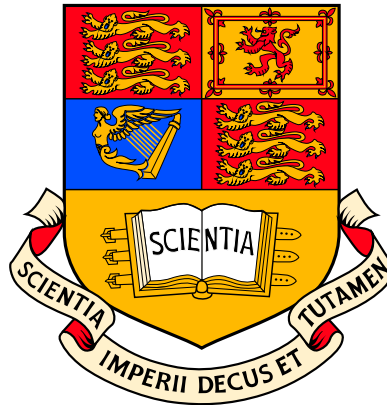


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Studies of Electron Acceleration Mechanisms in Relativistic Laser-Plasma Interactions

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Abstract

Laser-plasma interactions have many potential applications, such as medical treatments, x-ray generation, particle acceleration and inertial confinement fusion (ICF). In all of these applications, understanding how laser energy is absorbed by the material and converted into energetic electrons is very important. Therefore it is vital to enhance the understanding of how these energetic electrons are created and what mechanisms influence them.

This Thesis comprises experimental studies of electron acceleration mechanisms in laser-plasma interactions, as well as simulations relevant to these experiments. The experiments described were conducted at the Rutherford Appleton Laboratory utilising the VULCAN laser facility, and investigate laser interactions with both underdense and overdense plasmas.

In the underdense regime, the intensity dependence of the accelerated electrons has been studied experimentally, as well as the impact of the focusing geometry on the generation of hot electrons. For high intensities, experimental measurements show a scaling of the temperature of the electrons with a_0 . Density and f-number dependencies of the accelerated electrons are also observed.

The effect of laser polarisation and target thickness on the escaping electrons is studied for laser interactions with solid targets, or overdense plasmas. It was found that the effective temperature of the electrons depends on both the laser polarisation and the target thickness. The electron production from ultra-thin foils, and the effect of laser pre-pulse are also investigated.

The role of the author

In this section the role of the author in the work presented in this thesis is described. Experiments on the VULCAN laser system, involve a large team of people. The experimental data is usually collected by a team of experimental collaborators, with the technical assistance of the Central Laser Facility staff. The author was involved in the planning, setting up and running of the experimental diagnostics. In particular the electron spectrometers were run, and unless stated otherwise, the analysis of the experimental data was also conducted by the author. The author also helped to set up and run many of the other diagnostics on the experiments and in addition, was responsible for running two experiments in 2007 as the target area operator. The OSIRIS simulations presented in this thesis were all run and analysed by the author. The OSIRIS code was developed by the OSIRIS consortium (UCLA/USC/IST).

Declaration

I hereby certify that the material of this thesis, which I now submit for the award of Doctor of Philosophy, is entirely my own work unless otherwise cited or acknowledged within the body of text.

Signed,

Sabrina Roswitha Nagel

5th June 2009

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Chapter 1

Introduction

The advance in laser technology to short-pulse (≤ 1 ps) lasers has made it possible to study laser-plasma interactions at ultra-high intensities ($\geq 10^{18}$ Wcm $^{-2}$). This advance has been due to the implementation of chirped pulse amplification (CPA).

CPA was developed by Strickland and Mourou in the 1980s [1], and increased the peak power of lasers by orders of magnitude. In CPA the laser pulse is stretched in time before it is sent through the amplifiers, to avoid damage. After amplification, the pulse is re-compressed temporally. By focusing these high-energy, short-pulse laser pulses, intensities exceeding 10^{21} Wcm $^{-2}$ can be achieved [2]. This leads to a laser light pressure of $p_L = I/c = 3.33 \times 10^{11}$ bar and energy densities of around 1×10^{17} Jm $^{-3}$. At such high laser intensities, the normalised vector potential of the laser becomes greater than one, $a_0 \gg 1$. In the low intensity limit, a_0 is the quiver velocity, which is the ratio of the velocity at which an electron oscillates in the electric field of the laser to c . Hence as a_0 approaches one, the electron is driven to a velocity approaching the speed of light and the electron motion becomes relativistic. For an intensity of 10^{21} Wcm $^{-2}$, $a_0 = 29$, therefore the interaction is highly relativistic. Furthermore, because of the high laser intensities, any material the focused laser interacts with will be ionised and therefore form a plasma.

There are many applications of these high-intensity laser plasma interactions, including particle accelerators [3], inertial confinement fusion schemes [4] and also biological and medical technologies [5], described in more detail below. As a result

many very high intensity laser facilities have been built and have recently become available for research. Among them are the Omega EP laser at the University of Rochester, Titan at the Lawrence Livermore National Laboratory, the Texas Petawatt Laser at the University of Texas at Austin, the Gekko laser at the Osaka University, and the VULCAN PW laser at the Rutherford Appleton Laboratory, which was used for the experiments discussed in this thesis.

The generation of hot electrons in laser plasma interactions is important in many processes and applications. Therefore it is vital to study the acceleration of electrons in laser plasma interactions at ultra-high intensities over a variety of density regimes. This forms the topic of this thesis.

There are two major regimes of laser-plasma interactions, namely the study of laser interaction with underdense and overdense plasmas. In an overdense plasma the laser light cannot propagate, since the plasma density is larger than the “critical” density, where the electron plasma frequency equals the laser frequency ($\omega_p = \omega_0$). In contrast, an underdense plasma permits the propagation of the laser radiation. Some applications of laser plasma interactions are described in more detail below, first in the underdense and then in the overdense regime.

1.1 Laser interaction with underdense plasma

1.1.1 Plasma based particle acceleration

In conventional particle accelerators, the electric field used to accelerate charged particles is limited by electrical breakdown, to a maximum on the order of 50 MVm^{-1} [6]. Plasma based accelerators are unaffected by the problems associated with material breakdown at high field strength, since plasmas are already ionised. The use of laser excited plasma oscillations to accelerate electrons was first proposed by Tajima and Dawson [3]. Extensive research has shown that laser plasma accelerators are a promising alternative and/or supplement to conventional particle accelerators [7].

In a wakefield accelerator, particles gain energy from a longitudinal plasma wave. This plasma wave is set up as the ponderomotive force of the laser pulse pushes

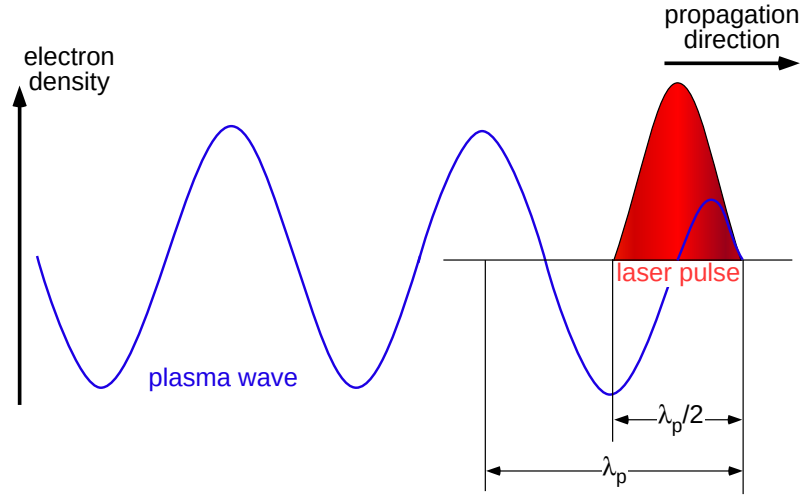


Figure 1.1: Plasma wave generation

electrons away from regions of high-intensity. When the pulse has passed, electrons rush back to their original position, since the ions remain in their original position, and a plasma wave is set up, see figure 1.1. These plasma waves reach the highest amplitude if the driving beam is shorter than the wavelength of the relativistic plasma wave ($\lambda_p = 2\pi c/\omega_p$). In fact the optimum pulse duration for a laser driven plasma wave is when the pulse length is half the relativistic plasma wavelength [3], as in the laser wakefield accelerator (LWFA). The frequency of the longitudinal electrostatic waves is $\omega_p [\text{s}^{-1}] = 56400\sqrt{n_e}$, where n_e is the electron density in cm^{-3} . These waves are limited by wavebreaking and can support electric fields on the order of $E [\text{V/m}] = 96\sqrt{n_e}$, which can be many GVm^{-1} even at moderate plasma densities.

If the laser pulse length is long compared to the plasma wavelength ($c\tau > \lambda_p$, where c is the speed of light and τ is the laser pulse duration), and the laser power P is larger than the critical power for relativistic self-focusing, $P \geq P_c [\text{GW}] = 17n_c/n_e$, the single pulse can break up into a train of short pulses. This break up is caused by a self-modulation instability. Experiments in this self-modulated laser wakefield acceleration (SMLWFA) regime demonstrated electron acceleration to above 100 MeV, but with a large energy spread [8–13].

At first, experiments in the LWFA regime, where the pulse length is shorter than the plasma wavelength, also showed broad electron spectra with energies up

to 200 MeV [14]. However, recently, high-energy, quasi-monoenergetic (low energy spread) electron bunches have been produced using the LWFA scheme [15–17]. It was demonstrated, that longer focal lengths ($\sigma_0 > \lambda_p$, where σ_0 is the waist of the focal spot), where the focal spot size is better matched to the plasma wavelength, improve the production of the electron beam, in terms of reproducibility, collimation and energy spread [18]. It has even been shown that it is possible to accelerate electrons to about 1 GeV over a distance of 10 mm using the laser wakefield accelerator [19,20].

The accelerated electrons, as they move in the plasma, can produce ultrashort-duration pulses of x-ray beams, by the synchrotron radiation mechanism [21]. This opens more possible applications for the plasma wave accelerators, such as γ -ray production, radioisotope production and possibly colliding-beam experiments.

1.1.2 Direct laser acceleration

The x-rays from these high-intensity laser plasma interactions can also be generated by electrons that are accelerated directly by the laser, in a scheme similar to the inverse free electron laser (IFEL) [22]. The principle of the IFEL is that electrons “wobble” in a magnetic field, caused by stationary magnets, at a frequency that depends on the electron energy. If an electromagnetic wave (i.e. laser) and a relativistic electron beam co-propagate through such a wiggler structure, the magnets produce a small transverse velocity (wiggling motion) in a direction parallel to the electric vector of the wave so that energy can be resonantly transferred between the particle and the wave. Because the acceleration rate depends on the electron velocity, the process is usually second-order. For this acceleration to happen, the magnetic field has to be designed, so that the electron’s wiggling motion and the electromagnetic wave are always in the same relative phase.

In intense laser plasma interactions, a comparable situation to the IFEL can be created, [23–25]. Ponderomotive expulsion of the electrons from the laser electric field sets up a channel, which creates a radial electric field. Electrons propagating in this channel will also set up an azimuthal magnetic field. These fields help to confine the fast electrons in the relativistic plasma channel, and the oscillations in the chan-

nel can lead to an effective energy gain, if the electrons are in the appropriate phase of the laser field. Therefore at ultra-high intensities, the effect of the electric and magnetic fields of the laser and within the plasma can lead to so-called direct laser acceleration (DLA). This direct laser acceleration regime leads to a broad energy spectrum and has been observed experimentally [26–28]. However, numerous dependencies on experimental parameters, such as the scaling of the electron temperature with laser intensity and electron density and the effect of the focusing geometry, are still to be determined.

The electrons from these interactions are also highly radiative, and generate x-rays, with up to 5% energy conversion, and x-rays with energies between 1 – 50 keV have been reported from laser-plasma interactions [21, 28]. The brightness of the x-ray source produced in such interactions is beginning to approach that of modern synchrotron sources, such as the Diamond light source [29], which has recently been made available for users and is part of the Rutherford Appleton Laboratory. Diamond produces x-rays with energies between 0.1 – 100 keV. Applications of synchrotron radiation are numerous and range from studying chemical reactions and compounds, analysing contaminants and understanding their release into the environment as well as their mobility and toxicity, investigating biological structures, studying fine structures, surfaces and magnetism through to medical diagnosis and treatment [30]. As third generation synchrotron facilities like Diamond are very expensive to both build and maintain, table top laser electron accelerators could bring the costs down, thus making comparable facilities available on a university scale.

1.2 Laser interactions with overdense plasma

Light cannot normally propagate into overdense plasma, above the critical density, n_c , defined as the density at which $\omega_0 = \omega_p$, where ω_0 is the laser frequency and ω_p is the plasma frequency. However laser energy can be absorbed by electrons at the critical surface. For high intensities, $I\lambda_\mu^2 > 10^{18} \text{ Wcm}^{-2}\mu\text{m}^2$, vacuum heating and $\mathbf{j} \times \mathbf{B}$ heating are the dominant absorption processes. These are described in more detail in chapter 2. The electrons that are accelerated leave the influence of the

laser as they travel into the target. When they leave the target at the rear, they set up electric sheath fields, and emit radiation as they transition into the vacuum.

In these interactions protons and ions can also be accelerated to energies of up to multiple MeV per nucleon. However, with the laser intensities available today, ions are too heavy to be accelerated directly, but are accelerated by secondary fields, generated by the plasma electrons.

1.2.1 Ion acceleration

As the laser-driven electrons try to leave the plasma, they set up an electric sheath field, which in turn accelerates the ions, mainly originating from hydrocarbon or water contaminants on the target surfaces. This mechanism for ion acceleration is the target normal sheath acceleration (TNSA), see figure 1.2(a), which so far is the most studied and efficient mechanism. The electric sheath fields are well above 1 TV/m. As the electric fields are varying in both time and space, the accelerated ions exhibit a large spread in energy and spatial distribution. However, monochromatic ion beams have been successfully generated by the use of reduced mass targets [31], microdot targets [32], thin layers of contaminant [33] and energy selection [34].

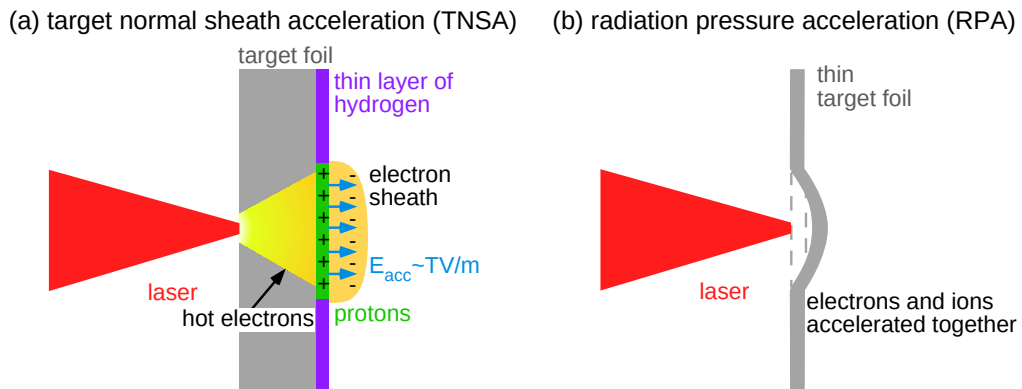


Figure 1.2: (a) schematic of the target normal sheath acceleration (TNSA) mechanism. (b) shows a sketch of the radiation pressure acceleration (RPA) mechanism.

Recently a scheme using ultra-thin targets in the laser piston regime [35, 36] has been proposed, which promises monoenergetic ion beams with high energies. This regime requires high-intensity as well as a high contrast ratio. Esirkepov *et al.* [35] demonstrated radiation pressure acceleration (RPA) numerically, at intensities $> 10^{23} \text{ Wcm}^{-2}$ using linear polarisation. Such intensities cannot be reached, with the lasers available today. Robinson *et al.* [36] therefore investigated the possibility of using circular polarisation to achieve a constant ponderomotive expulsion and to access the RPA regime with intensities around $10^{20} - 10^{21} \text{ Wcm}^{-2}$. In the RPA regime, the momentum of the laser is transferred to the particles more directly, see figure 1.2(b). Analytical models, based on momentum conservation, indicate that the final ion energy should scale as $\propto (I\tau/\sigma)^\alpha$, where I is the intensity, τ is the pulse duration, σ is the areal mass of the foil, and α is 2 for $v_{\text{final}} \ll c$ and $1/3$ for the ultra-relativistic limit.

Possible applications for these proton and ion beams are numerous, and are outlined only briefly. Ion beams can be used for cancer treatment, as they deposit their energy at a localised distance into the material, making the targeting of cancerous cells possible without damaging the healthy tissue [5, 37]. The challenge is to generate an ion beam with large enough flux at the required energy, using high repetition rate lasers, to decrease the size as well as the costs of treatment facilities, which would allow more people access to this treatment method.

The protons produced in laser plasma interactions, can themselves be used, to diagnose laser plasma interactions. This method is known as proton probing or proton radiography [38–40], and has been successfully used for time-resolved studies of the generation of electric and magnetic fields in the laser plasma interaction on a ps-timescale. These include the sheath electric fields at the rear of the target [41], the late-time ion structures in the wake of the high-intensity laser propagating through an underdense plasma [42] and self-generated magnetic fields around laser focal spots [43, 44].

1.2.2 Fusion

Another reason to study the acceleration of electrons and ions from overdense interactions is the hope of being able to harness fusion power to generate electricity. Two of the main energy resources, fossil fuels and fission, both have problems associated with them. Reserves of fossil fuels, such as oil, gas and coal are expected to be depleted in the near future and play a major role in global warming. On the other hand, the fission process produces radioactive material that requires storage for a significant time, and as the meltdown of the reactor in Chernobyl showed, can have disastrous effects on the environment and people. Therefore at the moment there is a big search for alternative energy resources to replace these, less favourable methods.

An alternative approach to produce energy is fusion, the process that powers the sun. Here small atoms are combined to form larger atoms while releasing energy in the process. If two light nuclei fuse, they will generally form a single nucleus with a slightly smaller mass than the sum of their original masses. The difference in mass, δm , is released in energy, according to Einstein's mass-energy relationship, $\delta \mathcal{E} = \delta m c^2$. Because of its low effective temperature threshold (10 keV), the DT reaction is the most promising fusion reaction to be achievable in the laboratory [45]. The equation below shows how the 17.6 MeV, released in the reaction, are distributed amongst the fusion products.



From this equation one can calculate an energy release of 340 GJ from 1 g of DT fuel. This is approximately the same as is released by burning 8 tonnes of oil. An additional advantage is that the terrestrial reserves of lithium, from which tritium can be obtained, and deuterium are abundant. To make fusion work however, the atoms must be heated to high temperatures and held at high pressure for long enough times for fusion to occur.

Two different routes to heat and confine the plasma in order to achieve fusion, are currently pursued. One is magnetic confinement fusion (MCF), where magnetic fields are used to confine the plasma. An international collaboration, ITER [46], aims

to demonstrate breakeven, using this technique with electron densities of 10^{14} cm^{-3} and confinement times of 3 – 4 s. Here breakeven refers to the energy produced in the fusion process being equal to the energy put in.

Inertial confinement fusion (ICF) is the second route, and uses strong irradiation, such as lasers, to compress and heat the hydrogen fuel to fusion temperatures, and uses the inertia of the fuel itself to confine it long enough for fusion to occur. The driver can be a laser, ion or electron beam, or a Z-pinch [47]. Large scale facilities built to investigate ICF, i.e. National Ignition Facility (NIF) [48] and Orion [49], are starting to come online. It is estimated that laser energies greater than a MJ are necessary to achieve ignition using the basic inertial confinement method. These high energies are necessary because the capsule has to be compressed at high temperatures for ignition to occur. NIF, which broke the 1 MJ barrier in March 2009, uses the indirect drive method where the pellet is compressed by the x-rays generated by the 192, nanosecond long pulses interacting with a solid target.

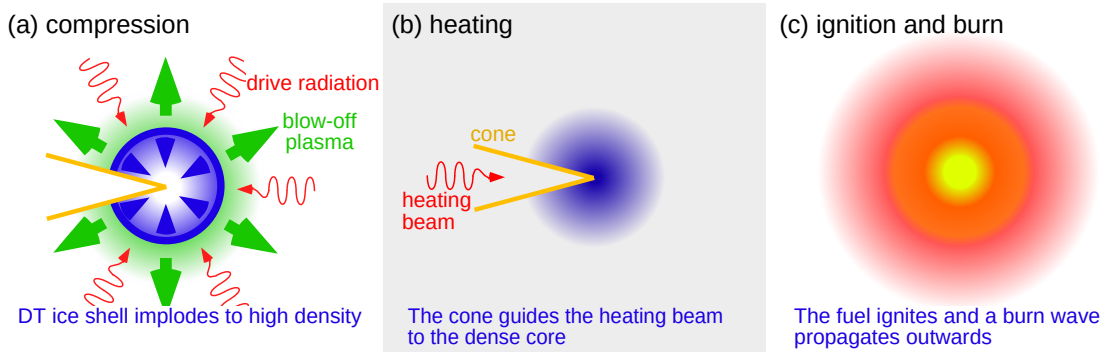


Figure 1.3: Schematic of the fast ignition approach to ICF. (a) shows the compression of the DT fuel capsule by radiation. (b) shows the guiding of the heating beam to the dense plasma by the cone. (c) shows the ignition of the dense fuel.

The fast ignitor (FI) concept was proposed by Tabak *et al.* [4], and is a variation of the ICF scheme, which separates the compression and heating processes. It offers the possibility of producing significantly higher gains with much less driver energy than the conventional approach. First, a capsule is imploded to form a high-density core, keeping the temperature low, see figure 1.3(a). The core is then ignited by a particle beam, created by the ignition laser. One problem of this scheme is deliv-

ering the particle beam to the core of the pellet, which has been “solved” by using a gold cone to keep the propagation path of the heating beam free of plasma, see figure 1.3(b). Using this scheme, the homogeneous compression of the fuel pellet and the effective heating of the compressed plasma, to temperatures of several hundred eV, using a high-intensity short-pulse laser, have been experimentally proven by Kodama *et al.* [50]. Li *et al.* [51] observed electron beam propagation along the target surface, indicating that the cone can guide electrons towards its tip.

The High Power laser Energy Research facility (HiPER) [52], which is an international project, is currently being designed to demonstrate the feasibility of laser driven fusion, utilising the FI scheme, as a future energy source.

1.2.3 Previous experiments

For all applications it is necessary to understand the generation of hot electrons, and therefore numerous experiments involving laser interactions with solid targets have been conducted. Here only some findings relevant to the thesis shall be reviewed.

The first experimental investigation into electron acceleration from the interaction of ultraintense ($I\lambda^2 > 10^{19} \text{ Wcm}^{-2}\mu\text{m}^2$), sub-picosecond laser pulses, with a solid target was described by Malka and Miquel [53]. Shooting a 30 μm thick CH target at normal incidence, with varying intensities (10^{18} to 10^{19} Wcm^{-2}), they found that the temperature of the hot electron tail scales like the ponderomotive potential of the laser ($T_{\text{eff}} \approx U_p$), as predicted by Wilks [54]. Later they also showed that femtosecond laser pulses are more efficient at generating energetic electrons than higher energy picosecond pulses. Electron temperatures up to 9.3 MeV for 6 μm thick plastic foils shot at normal incidence were observed, using a 1 J, 30 fs laser pulse, focused to an intensity of $2 \times 10^{19} \text{ Wcm}^{-2}$ [55]. However, their measurements for thicker targets were inconclusive, because of the acceptance angle of their diagnostic.

Cowan *et al.* reported the first observation of electrons accelerated up to 100 MeV from laser solid interactions, using the Petawatt laser at the Lawrence Livermore National Laboratory [56]. They attributed the high energy tail (20–100 MeV) to

self-focusing of the laser light or self-modulated laser-plasma acceleration in the underdense plasma region at the front of the target, created by the laser pre-pulse.

Concerning the scaling of the electron temperatures with the laser intensity, recent experiments [57] showed agreement with the Beg scaling [58], which suggests that the hot temperature of the escaping electrons scales as $(I\lambda^2)^{1/3}$. Haines *et al.* [59] introduced a model, which gives temperatures that closely agree with these experimental scalings [57, 58]. The model uses fully relativistic conservation equations and assumes 100% laser-light absorption.

The conversion efficiency of laser energy into electron energy is an important value, and the effect of target deformation on laser absorption has been investigated at $I\lambda^2 \approx 4 \times 10^{18} \text{ Wcm}^{-2}\mu\text{m}^2$. It was found that fractional absorption increases with growing deformation depth [60]. The deformation is due to a collisionless quasi steady bow shock [54, 61]. The conversion efficiency from laser energy into forward propagating electrons was measured to be 20–30% [62]. Nakatsutsumi *et al.* found that the absorption of laser energy into electrons entering the solid target was between 15 – 30% and tends to increase with laser intensity, which is in agreement with the above mentioned results [63]. Enhanced absorption at intensities above 10^{20} Wcm^{-2} was recently observed, and was found to be 60% for close to normal incidence and 80 – 90% for 45° incidence [64]. These absorption rates are just a little smaller than the ones predicted by the idealised model introduced in [59].

However, the conversion of laser energy into hot electrons is still not fully understood. This is highlighted by recent descriptions of collisionless absorption mechanisms [65, 66]. One model finds that preformed plasma is beneficial for high absorption, as the preformed plasma is compressed by the laser pulse [65].

After the electrons escape the influence of the laser, they travel through the target with a divergence angle, which is between $30^\circ - 60^\circ$, for intensities between $10^{19} - 10^{21} \text{ Wcm}^{-2}$. This was measured on different experiments and is summarised in [67]. Recent experiments showed the successful guiding of relativistic electron beams in solid targets, using resistively controlled magnetic fields [68].

Depending on the target thickness, the number of hot electrons within the target is different. It is expected to be higher for thinner targets, due to the effect of

recirculation [69, 70], which is described in more detail in chapter 2.

In relation to the RPA regime and ICF, recent simulations showed that for circular polarisation and intensities of around 10^{22} Wcm^{-2} the ion acceleration due to hole boring becomes efficient [71]. For the fast ignitor concept, this has the advantage that the pellet could be irradiated directly by the laser, and the fast ions from the outer target layers could lead to ignition of precompressed thermonuclear targets.

1.3 Thesis outline

In this thesis electron acceleration from intense laser interactions with both, underdense and overdense plasmas is studied. Experiments and numerical modeling were used to investigate the acceleration mechanisms. The work in this thesis covers areas relevant to electron acceleration, x-ray radiation generation, radiation pressure acceleration (RPA) and inertial confinement fusion (ICF). This thesis is organised in the following way.

Chapter 2 gives a short review of the relevant theory of laser-plasma interactions.

Chapter 3 describes the VULCAN laser system used to perform the experiments, and introduces the experimental methods used to obtain the data.

Chapter 4 presents the experiments where the focus was on electron acceleration from underdense plasmas. In particular it investigates interactions with intensities between $1.4 \times 10^{18} \text{ Wcm}^{-2}$ and $2 \times 10^{21} \text{ Wcm}^{-2}$ and densities ranging from $5 \times 10^{18} \text{ cm}^{-3}$ to $4 \times 10^{19} \text{ cm}^{-3}$. The chapter is split into two sections discussing the experiments with intensities above and below $1 \times 10^{20} \text{ Wcm}^{-2}$.

Chapters 5 and 6 on the other hand describe the electrons accelerated in laser interactions with overdense or solid targets.

Chapter 5 focuses on the interaction with ultra-thin, down to 50 nm thick, targets.

Chapter 6 investigates electron acceleration from 10s of μm thick targets. The dependence of the accelerated electrons on the laser polarisation, as well as the target thickness are discussed.

Chapter 7 summarises and concludes the work described in the thesis.

Chapter 2

Theoretical Review

A plasma is a quasineutral gas of charged and neutral particles which exhibits collective behaviour [72]. In gases, the particle motion is controlled by collisions between neutral particles. However, in a plasma, which has charged particles, inhomogeneous distributions (with non-uniform density) can give rise to electric fields. Motion of charges also creates currents, and therefore magnetic fields. These fields affect the motion of other charged particles far away. Collective behaviour means the motions depend not only on local conditions but on the state of the plasma.

Plasma in laser matter interactions is created by laser induced [73–77] or collisional ionisation [78]. Dependent on the laser intensities, different regimes of the laser induced ionisation can be accessed. The Keldysh parameter

$$\Gamma = \sqrt{\frac{\mathcal{E}_i}{2U_p}}$$

– where $\mathcal{E}_i$ is the ionisation energy and U_p is the ponderomotive potential (see equation (2.11)) – characterises these regimes. For $\Gamma \gg 1$ we are in the regime of sequential ionisation via virtual states (figure 2.1(a)), which means that a single electron has to absorb a large number of photons before it is released from its potential well. Electrons freed by multi-photon ionisation can knock out electrons from neutral particles, and therefore lead to avalanche breakdown in air, starting at intensities of about 10^{11} Wcm^{-2} . For $\Gamma \approx 1$ tunnelling ionisation can take place, where an electron can tunnel through the Coulomb barrier. Tunnelling ionisation (figure 2.1(b)) becomes the dominant process at intensities above $\approx 10^{14} \text{ Wcm}^{-2}$,

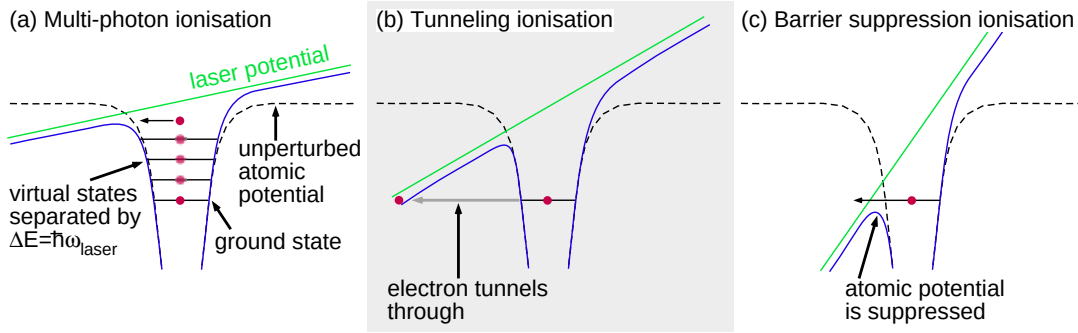


Figure 2.1: The three main Optical Field Ionisation (OFI) processes: (a) multi-photon ionisation, (b) tunneling ionisation and (c) barrier suppression ionisation.

when the electric field of the laser suppresses the Coulomb barrier. As the laser intensity reaches $\approx 10^{16} \text{ Wcm}^{-2}$ the Coulomb barrier is completely suppressed, and the regime of barrier suppression ionisation is accessed (figure 2.1(c)).

In high-intensity short pulse interactions, the plasma generation mechanism during the first few femtoseconds is field ionisation [77]. For solid targets, thinner than the laser wavelength, ultrafast ionisation can take place over the whole of the target volume [76]. However, determining the ionic charge-state distribution, of the plasma subjected to the rapidly changing conditions during short-pulse laser interactions is very complex. For high density, optically thick plasmas (where local thermal equilibrium (LTE) can be reached) the ionisation state can be determined by statistical means, and the relative ion populations are related by the Saha-Boltzmann equation [78]. However, short pulse lasers tend to produce optically thin plasmas, which span many orders of magnitude in density and temperature. For high atomic number elements, finding a solution to this non-LTE problem becomes very complicated.

2.1 Interaction with single atom and electron

2.1.1 Electron motion in an electromagnetic field

Plasma electrons are affected by the intense laser pulse, as they move in the electric and magnetic field of the laser. Let us first consider a single electron in an

electromagnetic field. A more detailed derivation of the electron trajectories in an electromagnetic field, following [79], can be found in appendix A.1. The equation of motion is given by the Lorentz force:

$$\frac{d\mathbf{p}}{dt} = q(\mathbf{E} + \mathbf{v} \times \mathbf{B}) \quad (2.1)$$

where $\mathbf{v}$ and $\mathbf{p} = \gamma m_e \mathbf{v}$ are the electron quiver velocity and quiver momentum respectively and $\gamma = (1 - v^2/c^2)^{-1/2} = (1 + p^2/m_e^2 c^2)^{1/2}$ is the relativistic Lorentz factor. Because of $|\mathbf{E}| = c|\mathbf{B}|$, only when the momentum of the electron, $\mathbf{p}$, becomes comparable to the rest mass of an electron, does the $\mathbf{v} \times \mathbf{B}$ force, that is due to the laser field, become significant. Therefore for an $a_0 < 1$, the electron motion will be primarily due to the electric field and will oscillate linearly along the electric field direction with the laser frequency. But in the nonlinear relativistic regime, the magnetic field component becomes non-negligible, and curves the electron trajectory towards the z -direction. As a result the electron has a longitudinal as well as transverse motion, with respect to the light propagation direction.

Using the vector identities $\mathbf{E} = -\frac{\partial \mathbf{A}}{\partial t} - \nabla \phi$ and $\mathbf{B} = \nabla \times \mathbf{A}$, and assuming a free electron in a plane electromagnetic wave propagating in the z -direction and polarised in the x -direction. The momentum in the z -direction of a particle with zero initial momentum can be found to be:

$$p_z = \frac{p_x^2}{2m_e c} = \frac{(qA_x)^2}{2m_e c} = m_e c \frac{a^2}{2} \quad (2.2)$$

where the normalised vector potential $a = qA/m_e c$ was defined. Equation (2.2) describes a parabolic trajectory in momentum space.

To derive the trajectory of the electron in space we need to integrate the equations of motion (see appendix A.1). We then obtain:

$$\begin{aligned} x(t) &= \frac{ca_0}{\omega} \sin(\omega t - kz) \\ z(t) &= \frac{ca_0^2}{4\omega} \left((\omega t - kz) + \frac{1}{2} \sin(2(\omega t - kz)) \right) \end{aligned} \quad (2.3)$$

Equation (2.3) indicates, that at low intensities, $a_0 \ll 1$, the motion is predominantly in the transverse direction. When $a_0 \gg 1$ the motion is predominantly longitudinal.

The two terms in the z -component of equation (2.3) indicate that the motion consists of a constant drift plus an oscillatory motion. If we move to a frame, that is moving at the drift velocity, the oscillatory motion of the electron describes a ‘figure-of-eight’ motion. Figure 2.2(a) shows plots of the trajectory of the electron and figure 2.2(b) is the ‘figure-of-eight’ trajectory in the drift frame of the electron. The drift

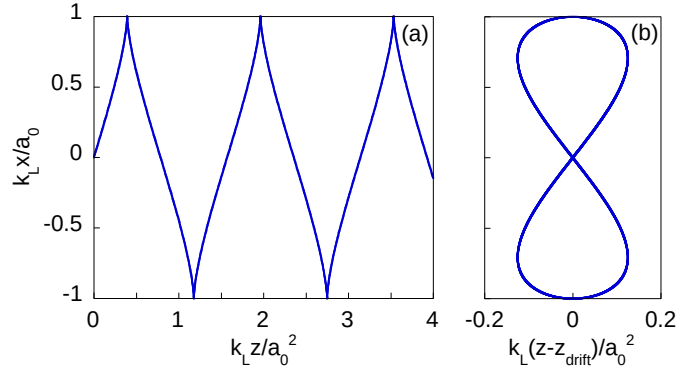


Figure 2.2: Plots for the case of a plane, linearly polarised electromagnetic wave showing (a) the trajectory of an electron and (b) the ‘figure-of-eight’ trajectory in the drift frame.

velocity, $\mathbf{v}_D$, as seen in the laboratory frame is found by the time averaging of the z component (see equation (A.18)), with $\langle \gamma \rangle = 1 + a_0^2/4$ for linear polarisation in vacuum, this becomes:

$$\mathbf{v}_D = \left\langle \frac{dz_{drift}}{dt} \right\rangle = \frac{a_0^2}{4 + a_0^2} c \mathbf{e}_z \quad (2.4)$$

For circular polarisation one gets [80]:

$$x = c \frac{a_0}{\omega} \sin(\omega\tau), \quad y = -c \frac{a_0}{\omega} \cos(\omega\tau), \quad z = ca_0^2 \tau \quad (2.5)$$

where τ also is: $\tau = t - z/c$. With $\langle \gamma \rangle = \sqrt{1 + a_0^2}$, for circular polarisation, the drift velocity becomes:

$$\mathbf{v}_D = \frac{a_0^2}{\sqrt{1 + a_0^2}} c \mathbf{e}_z \quad (2.6)$$

2.1.2 Laser propagation in a plasma

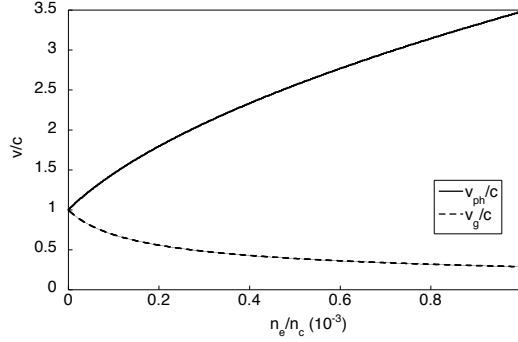


Figure 2.3: v_{ph}/c and v_g/c as a function of n_e/n_c .

The propagation of an electromagnetic wave through an underdense plasma is described by the plasma dispersion relation (for a derivation see appendix A.2) [81]:

$$\omega^2 = \omega_p^2 + c^2 k^2 \quad (2.7)$$

ω_p in equation (2.7) is the non-relativistic plasma frequency defined as:

$$\omega_p = \sqrt{\frac{n_e e^2}{m_e \epsilon_0}} \quad (2.8)$$

For $a_0 \rightarrow 1$, m_e can be replaced with $\langle \gamma \rangle m_e$, to get the relativistic electron plasma frequency.

Figure 2.3 shows the phase velocity ($v_{ph} = \omega/k = c(1 - n_e/n_c)^{-1/2}$) and the group velocity ($v_g = \partial\omega/\partial k = c(1 - n_e/n_c)^{1/2}$) of the laser, normalised to c , in the plasma, as a function of n_e/n_c . The group velocity decreases with increasing density, whereas the phase velocity increases.

An electromagnetic wave is only able to propagate if its frequency $\omega_0 > \omega_p$. $\omega_0 = \sqrt{\frac{n_c e^2}{m_e \epsilon_0}}$, where n_c is the critical density above which the laser is unable to propagate.

$$n_c = \frac{m_e \epsilon_0 \langle \gamma \rangle \omega_0^2}{e^2} = 1.12 \times 10^{21} \frac{\langle \gamma \rangle}{\lambda_\mu^2} [\text{cm}^{-3}] \quad (2.9)$$

Plasmas with densities below the critical density, $n_e < n_c$, are called underdense, whereas plasmas with densities above the critical density, $n_e > n_c$, are referred to as overdense plasmas.

2.1.3 Ponderomotive force

The time average of the oscillatory motion of the electron in the laser field gives the ponderomotive force (a derivation can be found in appendix A.3). The ponderomotive force can be written as [72, 81, 82]:

$$\mathbf{F}_p = -\frac{mc^2}{2\langle\gamma\rangle}\nabla\langle\mathbf{a}^2\rangle = -\frac{mc^2}{2\langle\gamma\rangle}\nabla(\langle\gamma\rangle^2 - 1) = -mc^2\nabla\langle\gamma\rangle \quad (2.10)$$

where $\langle\rangle$ denotes time-averaging over the fast laser cycle. For linearly polarised light, $\langle\mathbf{A}^2\rangle = \frac{1}{2}A_0^2$ and $\langle\mathbf{a}^2\rangle = \frac{1}{2}a_0^2$, and for circularly polarised light, $\langle\mathbf{A}^2\rangle = A_0^2$ and $\langle\mathbf{a}^2\rangle = a_0^2$. The Lorentz factor can be written as $\langle\gamma\rangle = \sqrt{1 + \langle\mathbf{a}^2\rangle}$ ^{*}. According to equation (2.10), an electron is expelled from high-intensity regions of the laser focus along the gradient of the laser-intensity distribution.

From equation (2.10) the ponderomotive potential can be found

$$U_p = \frac{mc^2}{2\langle\gamma\rangle}\langle\mathbf{a}^2\rangle = mc^2(\langle\gamma\rangle - 1) \quad (2.11)$$

which is the energy an electron gains as it quivers in the field of the laser.

2.1.4 Debye length

In the interactions described, the quasi-neutrality of the plasma is often disturbed and large electrostatic fields are generated. Because of their smaller mass-to-charge ratio, the electrons usually respond to these fields and move to compensate for them. Assuming that the ion temperature is much less than the electron temperature, the Debye length is the distance over which the electric field is screened.

$$\lambda_D = \sqrt{\frac{\epsilon_0 k_B T_e}{e^2 n_0}} \quad (2.12)$$

^{*}This is due to $\gamma = \sqrt{1 + p_{\parallel}^2 + p_{\perp}^2}$. But for the propagation in a plasma the drift velocity is not important, due to the shielding by plasma electrons, and can be neglected. Therefore $\langle\gamma\rangle$ becomes $\langle\gamma\rangle = \sqrt{1 + a_0^2/2}$ for linear polarisation and $\langle\gamma\rangle = \sqrt{1 + a_0^2}$ for circular polarisation.

2.2 Laser interactions with underdense plasma

2.2.1 Relativistic self-focusing and optical guiding of laser pulses

In this section we are going to look at the refractive index and how it leads to self-focusing, following Esarey *et al.* [6].

The ponderomotive force expels the electrons from the region of highest intensity

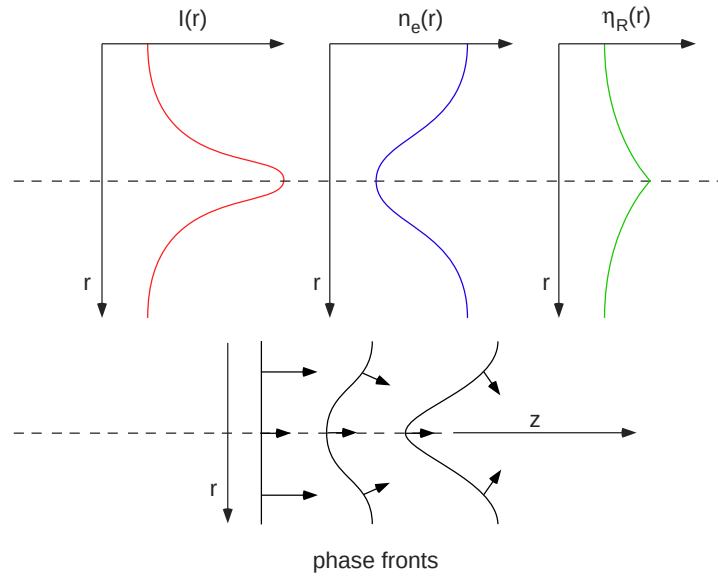


Figure 2.4: Diagram of the radial intensity, electron density and refractive index (top) as well as the evolution of the phase front.

and therefore lowers the electron density, n_e , in the central region. The refractive index $\eta_R(r)$ exhibits a maximum on axis, $\frac{\partial \eta_R}{\partial r} < 0$, since $\eta_R \approx \frac{ck}{\omega_0}$, $\frac{\partial \eta_R}{\partial r} < 0$ implies that the phase velocity along the propagation axis is less than it is off axis. This causes the laser phase front to curve such that the beam focuses towards the axis. For small amplitude electromagnetic waves, $n_e = n_0$, where n_0 is the ambient electron density, and

$$\eta_R = \frac{ck}{\omega_0} = \left(1 - \frac{\omega_p^2}{\omega_0^2}\right)^{1/2} \quad (2.13)$$

For large amplitude electromagnetic waves variations in the electron density and mass will occur: $\omega_p^2 \rightarrow \left(\frac{\omega_p^2}{\gamma}\right) \frac{n_e}{n_0}$, assuming $\frac{\omega_p^2}{\omega_0^2} \ll 1$, $\eta_R(r)$ can be modified through

$\gamma(r)$ and $n_e(r)$.

$$\eta_R \approx 1 - \frac{\omega_p^2}{2\omega_0^2} \frac{n_e(r)}{n_0\gamma(r)} \quad (2.14)$$

Since the leading motion of the electrons in the laser field is the quiver motion, $\mathbf{p}_\perp = m\mathbf{c}\mathbf{a}$, so $\gamma \approx \gamma_\perp = (1 + a^2)^{1/2}$ (for $a < 1$). A laser intensity profile peaked on axis, $\frac{\partial a^2}{\partial r} < 0$, leads to $\frac{\partial \eta_R}{\partial r} < 0$ and the possibility of guiding (i.e. relativistic self-focusing). The density profile can have contributions from a pre-formed density channel $\Delta n_p \sim \Delta n r^2 / r_0^2$ or a plasma wave $\delta n_e \sim \delta n_0(r) \cos(k_p \zeta)$, i.e. $n_e = n_0 + \Delta n_p + \delta n$. A radial density profile, which has a minimum on axis (i.e. a channel) implies $\partial \eta_R / \partial r < 0$. In the limits $a^2 \ll 1$, $|\Delta n_p / n_0| \ll 1$ and $|\delta n / n_0| \ll 1$, the refractive index can be written as [6]:

$$\eta_R \simeq 1 - \frac{\omega_p^2}{2\omega_0^2} \left(1 - \frac{a^2}{2} + \frac{\Delta n_p}{n_0} + \frac{\delta n}{n_0} \right) \quad (2.15)$$

In the above expression the $a^2/2$ term (which derives from the approximation: $(1 + a^2)^{-1/2} \simeq 1 - a^2/2$ in the $a^2 \ll 1$ limit) is responsible for relativistic optical guiding, the $\Delta n_p / n_0$ term is responsible for preformed density channel guiding, and the $\delta n / n_0$ term is responsible for self-channeling, plasma wave guiding and self-modulation of long laser pulses. The self-focusing of the beam competes against diffraction, the natural tendency for the beam to defocus. Relativistic self-focusing can overcome diffraction when the laser pulse power reaches a critical power $P \geq P_c$ [83–85]

$$P_c = \frac{8\pi\epsilon_0 m_e^2 c^5 \omega_0^2}{e^2 \omega_p^2} = 17 \frac{n_c}{n_e} \quad [\text{GW}] \quad (2.16)$$

A derivation of the critical power can be found in appendix A.4.

Relativistic self-guiding has been experimentally demonstrated, amongst others in [86–90].

2.2.2 Instabilities in laser produced plasmas

Modulation instabilities arise when the laser pulse wave packet is modulated by changes in the refractive index. Since the laser envelope is described by $\langle a^2 \rangle$, a modulation of this vector potential results in a change to the ponderomotive force of the laser.

One instability that affects the laser propagation is filamentation, in which the beam breaks up. Depending on the laser power, different mechanisms can lead to filamentation. For $P < P_c$ filamentation can occur due to side scattering instabilities, such as the ionisation induced scattering instability [91] and stimulated Raman side scattering [92]. For powers much greater than P_c , the non-linear evolution of the self-focusing towards self-guiding or filamentation depends on several factors, among them the value of P/P_c and the quality of the incident beam, as variations in the beam profile or phase front can affect filamentation [93–95].

Modulations (transverse) in the laser intensity, can lead to self-focusing of parts of the beam (due to either ponderomotive or relativistic effects). This results in a lower plasma density in the filament, and therefore even tighter focusing, leading to an unstable feedback loop where a tighter focus leads to higher intensity and more density depletion. Filamentation has been observed in both experiments [18, 96–98] and simulations [94, 99]. In laser-plasma electron acceleration experiments filamentation usually limits the energy of the accelerated electrons. This is due the lower intensity in each filament, compared to the laser intensity the beam would have without the break up.

The most important instabilities commonly encountered are the Stimulated Raman Scattering (SRS) and the Stimulated Brillouin Scattering (SBS) instability, which are parametric instabilities. Both instabilities involve the decay of an incident photon into a scattered photon plus an electron plasma wave for the SRS instabilities and an ion acoustic wave for the SBS instability when the frequency and wave number matching conditions are satisfied. In the SRS case, the matching condition is [81]:

$$\omega_0 = \omega_s + \omega_p \quad (2.17)$$

$$\mathbf{k}_0 = \mathbf{k}_s + \mathbf{k}_p \quad (2.18)$$

where ω_0 , ω_s and ω_p are the incident light, scattered light and plasma frequencies and $\mathbf{k}_0$, $\mathbf{k}_s$ and $\mathbf{k}_p$ are the incident light, scattered light and plasma wave-numbers respectively. Since the minimum frequency of a light wave propagating in a plasma is ω_p , this instability requires that $\omega_0 \gtrsim 2\omega_p$, which corresponds to the density $n \lesssim \frac{n_c}{4}$.

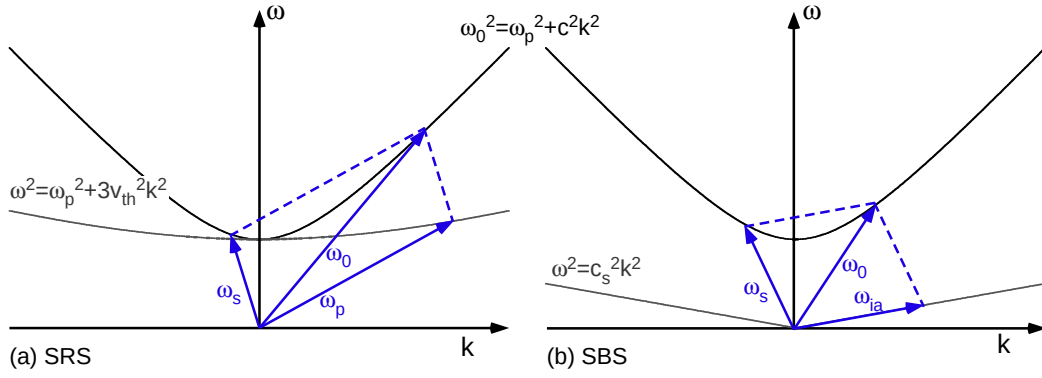


Figure 2.5: Sketch of the (a) stimulated-Raman-scattering and (b) stimulated-Brillouin-scattering mechanisms. ω_{ia} denotes the frequency of the ion acoustic wave.

The Raman instability leads to modulations in the spectrum of the laser light as it is passing through the plasma. It is possible for a forward Raman scatter mode to undergo multiple scattering, sometimes referred to as cascading [6]. As a consequence, the main frequency of the laser pulse will have additional satellites which have the frequencies given by:

$$\omega_s = \omega_0 \pm n\omega_p \quad (2.19)$$

where n is an integer. These satellites can also be due to scattering off a nonlinear plasma wave [100]. The positions of these peaks give the plasma frequency, which in turn gives the plasma density, $n_e = \frac{\epsilon_0 m_e \omega_p^2}{e^2}$. Therefore the Raman instability offers a way to measure the electron density in underdense plasma experiments [101].

2.2.3 Electron acceleration mechanisms from underdense plasmas

There are several mechanisms through which electrons can be accelerated from laser interactions with underdense plasmas. They can be divided into two regimes, the plasma wave accelerator regime and the direct laser acceleration regime. The first part of this section discusses the energy gain from plasma waves, and in section 2.2.3.7 the different mechanisms leading to direct laser acceleration are explained.

2.2.3.1 Plasma waves

The acceleration of electrons by plasma waves generated by intense laser plasma interactions was first suggested in 1979 by Tajima and Dawson [3]. The fundamental principle being that the ponderomotive force of an intense electromagnetic wave can set up a plasma wave which travels with a relativistic phase velocity. Due to their large longitudinal electric fields, relativistic plasma waves can accelerate charged particles.

The plasma wave is set up as the ponderomotive force pushes electrons away from the regions of high intensity. Due to their higher mass to charge ratio the ions remain within the laser pulse, so a space charge field is set up which pulls the electrons back towards their original position. This leads to oscillations at the plasma frequency ω_p (for a derivation see appendix A.5). The laser travels through the plasma at the group velocity (see section 2.1.2). As it propagates it continually sets up a plasma wave at each point along the path. Each oscillating region is essentially independent. However, the phase difference between a region where the oscillation started at $t = 0$ and that where the oscillation started at $t = \Delta t$ is $\Delta\Phi = v_g\Delta t$. Therefore the laser drives a plasma wave with a relativistic phase velocity equal to the group velocity of the laser pulse and the plasma wave has a wavelength of $\lambda_p = \frac{2\pi v_g}{\omega_p}$.

2.2.3.2 Wavebreaking

Wavebreaking occurs, because some of the plasma electrons undergo such large oscillations that the returning force due to the plasma wave is no longer large enough to make them continue their longitudinal oscillation (i.e. when the amplitude of the displacement of electrons in the wave approaches the plasma wavelength). They can then become trapped and accelerated. Therefore wavebreaking is a way to inject electrons into the plasma wave to be accelerated [101–103]. Once electrons have been trapped by the wave they are effectively “surfing” the wave and can gain energy from it.

The breaking of the plasma wave limits the electric field strength achievable, and therefore the energy gain possible. Using Poisson’s equation ($\nabla^2\phi = -\frac{\rho}{\epsilon_0}$)

and assuming the oscillations are harmonic, we can derive a maximum electric field associated with this:

$$\nabla \cdot \mathbf{E} = -e\delta n/\epsilon_0 \quad (2.20)$$

$$|i\mathbf{k}_p \cdot \mathbf{E}_0| = -en_0/\epsilon_0 \quad (2.21)$$

If we assume that the phase velocity of the plasma wave approaches the speed of light so that $k_p = \omega_p/c$ we recover:

$$E_0 = m_e c \omega_p / e \quad (2.22)$$

The above expression is a non-relativistic estimate. A fully relativistic one-dimensional calculation [104–106] is extensive, and here only the solution for the maximum electric field in the cold wavebreaking limit shall be presented.

$$E_{WB} = \sqrt{2(\gamma_{ph} - 1)} \frac{m_e c \omega_p}{e} = E_0 \sqrt{2(\gamma_{ph} - 1)} \quad (2.23)$$

where $\gamma_{ph} = (1 - \beta_{ph}^2)^{-1/2}$ and $\beta_{ph} = v_{ph}/c$. Including thermal effects reduces the threshold for wavebreaking to occur [107]. When the threshold is exceeded, the wave structure is destroyed. However, if there is a constant driver, i.e. a laser pulse, the wave is maintained, but some of the electrons will no longer be coherent with the plasma wave, and will therefore experience an accelerating field. This process is called trapping.

2.2.3.3 Trapping

Electrons can only be trapped and accelerated by a plasma wave if they are travelling at some initial velocity close to the phase velocity of the wave. Particle trapping and wavebreaking are closely linked. During wavebreaking wave energy is transferred directly to the particles.

However, the wavebreaking potential is the potential that is required to turn around electrons that are traveling backwards at the wave phase velocity. Therefore electrons that are not part of the wave structure, and have a velocity close to the wake phase velocity (which is close to the group velocity of the laser pulse in the plasma, i.e. $v_g = c\sqrt{1 - \frac{\omega_p^2}{\omega_0^2}}$), can be trapped at a lower potential than the wavebreaking potential [6, 10, 108].

2.2.3.4 Laser Wakefield Acceleration

In the laser wakefield accelerator (LWFA), a single short pulse ($\lesssim 1$ ps) ultra-high intensity ($\gtrsim 10^{18}$ Wcm $^{-2}$) laser pulse drives a plasma wave. The wakefield is driven most efficiently when the laser pulse length $c\tau$ is approximately the plasma wavelength $\lambda_p = 2\pi c/\omega_p$, i.e. $c\tau \simeq \lambda_p$ [6]. In order to accelerate electrons they have to be trapped by the plasma wave (chapter 2.2.3.6) and can then be accelerated by the large longitudinal field of the plasma wave (chapter 2.2.3.1).

As an intense laser pulse propagates through an underdense plasma, $\lambda^2/\lambda_p^2 \ll 1$, the ponderomotive force associated with the laser pulse envelope (equation (2.10)) expels electrons from the region of the laser pulse. If $L_z \sim \lambda_p$, where L_z is the axial gradient of the pulse profile, the ponderomotive force excites a large amplitude plasma wave (wakefield). This plasma wave has a phase velocity approximately equal to the laser group velocity [85, 104, 109, 110].

For an optimised flat-top pump pulse, that is linearly polarised, the accelerating field is given by $E_{max}/E_0 = a_0^2/2(1 + a_0^2/2)^{-1/2}$, where $E_0 = m_e c \omega_p / e$ [111]. The wakefield amplitude reaches its maximum, when the laser pulse length $L \simeq \lambda_{Np}/2$ where $\lambda_{Np} = \frac{2}{\pi} \frac{E_{max}}{E_0} \lambda_p$ is the nonlinear plasma wavelength [6]. The wakefield electric field can be increased by going to high densities and shorter pulse lengths, since $E_{max} \propto \lambda_p^{-1}$. However, at high densities, the laser group velocity is reduced and electron phase detuning can limit the energy gain (see section 2.2.3.6).

The generation of relativistic electron beams with low divergence and small energy spread by this mechanism has been experimentally observed [15–17].

2.2.3.5 Self-Modulated Laser Wakefield Acceleration

The self-modulated laser wakefield accelerator (SMLWFA) uses a single short ($\lesssim 1$ ps) ultrahigh intensity ($\gtrsim 10^{18}$ Wcm $^{-2}$) laser pulse, as in the LWFA. The SMLWFA, however operates at higher densities than the standard LWFA, so that the laser pulse length is long compared to the plasma wavelength $c\tau > \lambda_p$, and the laser power P is larger than the critical power P_c for relativistic guiding, $P \geq P_c$. In this high-density regime, the laser pulse undergoes a self-modulation instability (see section

2.2.2) which causes it to become axially modulated at the plasma period. Associated with the modulated pulse structure is a large amplitude, resonantly driven plasma wave. The self-modulation instability is a three-dimensional version of the forward Raman instability.

A single long pulse laser can, under appropriate conditions, break up into a train of short pulses. Each of these pulses has a width on the order of λ_p . Associated with the breakup of the long pulse is a large amplitude wave. In other words, a long pulse $\tau \gg 1/\omega_p$ in a plasma will couple to plasma waves via the Raman forward scattering instability. These waves gain in amplitude until they are sufficient for a combination of longitudinal bunching of the photons and transverse focusing of the pulse to cause a train of pulses modulated at the plasma frequency.

Electrons have been accelerated in excess of 100 MeV in SMLWFA experiments [8–13].

2.2.3.6 Electron energy gain from plasma waves

In section 2.2.3.2 we saw that wavebreaking represents a limit to the electric field strength achievable in the plasma wave. Since the maximum energy an electron can gain from a plasma wave will be determined by the product of the electric field strength of the plasma wave and the acceleration distance, the wavebreaking limit represents a limit to the energy gain of the electron. One other limiting factor in laser plasma acceleration is the length over which it is possible to accelerate electrons. This is dependent on several limiting factors. First of all, the interaction region can be limited by the characteristic diffraction of the laser, i.e. Rayleigh length $z_{RG} = \pi w_0^2/\lambda$. For a non-gaussian beam, this is, $z_R = \frac{z_{RG}}{M^2}$, where the M^2 -factor is a parameter for quantifying the quality of the laser (also see appendix A.6) [112]. This limitation, due to diffraction, can be overcome to some extent either by the use of channel guiding [113–116] or by using the self-guiding [20, 88] that occurs due to relativistic and plasma wave self-focusing. The mechanism of self-focusing is described in section 2.2.1.

Dephasing length

Since relativistic plasma waves effectively travel with the group velocity of the laser pulse that generated them, there is a mismatch between the speed of the wave and the speed of the electrons. To reach the maximum energy an electron will need to move from the potential peak to a potential trough, i.e. half a plasma wavelength. Then it will overtake the wave and begin to be decelerated by the second half of the plasma wave. If the electron has a velocity $v = \beta_e c$ and the plasma wave has a phase velocity $v_\phi = \beta_\phi c$, then the time it takes for this to occur is $t_d = \frac{\lambda_p}{2c(\beta_e - \beta_\phi)}$. The dephasing length is the distance the electron travels in this time. We can assume that $\beta_e \rightarrow 1$ and that the phase velocity of the plasma wave is given by the group velocity of the laser pulse, i.e. $\beta_\phi = \frac{v_g}{c} = \left(1 - \frac{n_e}{n_c}\right)^{\frac{1}{2}} \approx 1 - \frac{1}{2} \frac{n_e}{n_c}$, since the plasma is underdense. This reduces to:

$$L_d = \frac{n_c}{n_e} \lambda_p = \gamma^2 \lambda_p \quad (2.24)$$

In [111] the linear and non-linear cases in 1D are discerned:

$$L_d = \frac{n_c}{n_e} \lambda_p \begin{cases} 1, & \text{for } a_0^2 \ll 1 \\ \sqrt{2} a_0 / \pi, & \text{for } a_0^2 \gg 1 \end{cases} \quad (2.25)$$

For $a_0^2 \gg 1$ the dephasing length becomes longer, due to non-linear lengthening of $\lambda_p \rightarrow \lambda_{Np}$ (see section 2.2.3.4).

Laser depletion length

Another limiting factor is the depletion of the laser pulse energy. The plasma wave generated has an energy associated with it: $W_{plasma} = U_{plasma} V_{plasma} = \frac{1}{2} \epsilon_0 E_z^2 \pi w_0^2 L$, where U_{plasma} is the energy density of the plasma wave, V_{plasma} is the plasma volume and E_z is the electric field of the plasma wave. And the total energy in the laser is: $W_L = U_L V_L = \epsilon_0 E_L^2 \pi w_0^2 c \tau$. The length over which the laser loses energy can be estimated by $W_L = W_{plasma}$. For a circularly polarised wave, and $a_0 \ll 1$, the plasma wave electric field is given by $E_z = a_0^2 m_e c \omega_p / e^\dagger$. This occurs if the laser pulse length is such that $c \tau = \lambda_p / 2$. The laser pulse electric field can be expressed in terms of

[†]Since $E_{max} = a_0^2 (1 + a_0^2)^{-1/2} E_0$ for circular polarisation, which is $\approx a_0^2 E_0$ for $a_0 \ll 1$ [6,111].

the vector potential through $E_L = a_0 m_e c \omega_0 / e$. Combining these equations provides us with the pump depletion length in the linear regime for $a_0 \ll 1$:

$$L_{pd} = \frac{n_c}{n_e} \frac{\lambda_p}{a_0^2} = \gamma^2 \frac{\lambda_p}{a_0^2} \quad (2.26)$$

Equation (2.26) implies that increasing the laser energy for a given spot size and pulse duration does not increase the depletion length, on the contrary it decreases it.

In [111] the linear and non-linear cases in 1D are discerned.

$$L_{pd} = \frac{n_c}{n_e} \lambda_p \begin{cases} 4/a_0^2, & \text{for } a_0^2 \ll 1 \\ 2\sqrt{2}a_0/\pi, & \text{for } a_0^2 \gg 1 \end{cases} \quad (2.27)$$

For $a_0^2 \gg 1$ the depletion length also becomes longer, due to non-linear lengthening of $\lambda_p \rightarrow \lambda_{Np}$, where $\lambda_{Np} = \lambda_p \sqrt{2}a_0/\pi$ for linear polarisation and flat-top pulse shape (see section 2.2.3.4).

Limits of energy gain

In order to obtain the energy gain of an electron in a plasma wave, we have to integrate the longitudinal electric field of the plasma wave over the distance that the electron covers while it is in the wave: $W_{max} = e \int E_z dz$. Assuming that the amplitude of the plasma wave is constant (in reality it may be changing with time) at the maximum amplitude, $E_z = a_0^2 E_0$ then the maximum energy gain is given by:

$$W_{max} = 2 m_e c^2 a_0^2 \frac{n_c}{n_e} \quad (2.28)$$

In [6] the linear and non-linear cases for linear polarisation are discerned.

$$W_{max} = 2 m_e c^2 \frac{n_c}{n_e} \begin{cases} a_0^2/2, & \text{for } a_0^2 \ll 1 \\ a_0/\sqrt{2}, & \text{for } a_0^2 \gg 1 \end{cases} \quad (2.29)$$

In [117] Lu *et al.* compare formulas and scalings from linear theory, 1D non-linear theory and 3D theory.

2.2.3.7 Direct Laser Acceleration

At relativistic intensities, the electron moves mainly in the direction of the laser propagation (see section 2.1.1), but it does not exceed the ponderomotive energy, and falls back to its initial energy after the laser pulse has overtaken it. Therefore there is no net energy transfer to the electron (Lawson–Woodward criterion [118]). This behaviour reflects the fact that photons cannot be absorbed by electrons unless there is momentum exchange with third partners. However, in the relativistic regime ($a_0 \gg 1$), at ultra-high intensities the effect of the laser magnetic field and electric and magnetic fields within the plasma can lead to an acceleration in the longitudinal direction of the laser propagation. This direct laser acceleration regime has been observed experimentally in [26–28]. The three mechanisms that can lead to an energy gain of the electron directly from the laser, are described here.

A focused laser will generally produce a channel in the plasma through ponderomotive expulsion of electrons from the beam axis. Accompanying the pulse inside this channel, a current of fast electrons with a velocity close to the speed of light ($v \sim c$) may propagate, driven by wakefield acceleration or by one of the DLA mechanisms. In such a case there will exist a radial electric field E_s and an azimuthal magnetic field B_s within the channel that can be regarded as static over the duration of the pulse.

- Betatron resonance

Pukhov *et al.* [23] first described this mechanism of direct laser energy coupling to hot plasma electrons. It requires strong self-generated static electric and/or magnetic fields. These are generated by plasma electron depletion under the ponderomotive force, which creates a radial electric field, as well as by the stream of accelerating electrons, which creates an azimuthal magnetic field. These fields are both focusing for electrons traveling in the laser direction and so any such electrons will oscillate in the channel, called betatron oscillations. Energy transfer is possible when the betatron frequency ($\omega_\beta = 1 - v_z/v_{ph}$) is close to the laser frequency, as seen by the relativistic electron, so the Doppler shifted laser frequency $\omega_d = \omega_L \left(1 - \frac{v_z}{v_{ph}}\right)$. The derivations for the betatron

frequency as well as the energy at which the resonance occurs,

$$\gamma_e \simeq 4 \frac{n_c^2}{n_e^2} \left(1 + \frac{a_0^2}{4} \right)^2 \nabla \mathbf{F},$$

– where $\mathbf{F}$ is the focusing force, due to the combination of magnetic and electric fields – can be found in appendix A.7. It can be found that the energy an electron has to reach, before it can gain energy from the resonance, depends on background plasma density as well as laser intensity.

The electron spectra that Pukhov *et al.* obtained from their 3D–PIC simulations, show that the electrons are accelerated to above the ponderomotive energy and have a Boltzmann–like form with effective temperatures T_{eff} growing with laser intensity as $T_{\text{eff}} \propto \sqrt{I}$ [23]. They also undertook simulations comparing linear and circular polarisation of the laser. For the linearly polarised laser beam the transverse velocity of the resonant electrons was found to oscillate with the laser period, while the longitudinal velocity oscillated at twice per laser period, leading to a bunching of the electrons at $2\omega_0$. In the case of a circularly polarised beam, this bunching was absent.

The best acceleration is achieved when the betatron oscillations and the laser electric field are inversely phased, so that as the electron reaches the channel boundary, and its turning point, the electric field vanishes. Further studies of this mechanism by Tsakiris *et al.* [25] showed that the energy gain is large for larger $\eta_R^\dagger$, as phase matching conditions for these electrons can be satisfied more easily.

As the electron gains energy, ω_β decreases, detuning the electron, so that the resonance condition is not met. In the case of on–axis born electrons, only those accelerated electrons retain their energy which move ahead of the peak of the pulse. The energy of the ones that are run over by the pulse is largely given back to the pulse. Of course if the length of the channel is truncated before that, then the electrons can retain the energy.

[†] η_R is the refractive index, and is important, because if the initial velocity of the electron in laser propagation direction $v_{z0} < v_g = \eta_R c$, then the maximum velocity the electrons can gain on axis is also $< v_g$.

- B-loop mechanism

An electron quivering inside a relativistically intense laser pulse can be diverted in the transverse direction with a non-zero energy gain (ponderomotive expulsion). As it makes a loop in the surrounding azimuthal B -field it can be reinjected back into the laser pulse. This magnetic field must have a focusing polarity, corresponding to a counter-clockwise direction when looking along the laser, and can be set up i.e. by pre-accelerated electrons. It also has to have a sufficient extent r_B in the radial direction, and has to be strong enough to send back escaping electrons. The condition for re-injection of relativistic electrons with $p_{\parallel} \gg p_{\perp}$ was estimated to be [119]:

$$B r_B > \frac{a_0^2}{2\gamma} B_0 \lambda_0,$$

where $B_0 = m_e \omega_0 / e = 107$ MG, for a laser pulse with $\lambda_0 = 1 \mu\text{m}$ and $\gamma = \gamma_0 \{1 + a_0^2 / [\gamma_0^2 (\beta_{ph}^2 - 1)]\}$ with $\beta_{ph} = v_{ph} / c$ and $\gamma_0 = \sqrt{1 + p_0^2}$. A schematic of this mechanism can be seen in figure 2.6.

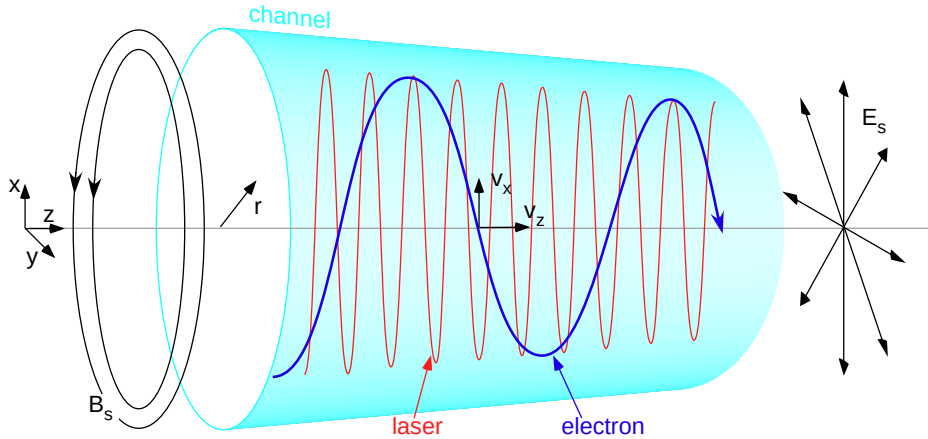


Figure 2.6: Direct laser acceleration (B-loop mechanism)

As the electron quivers in the electric field of the laser, the radial velocity and distance to the centre of the channel increase, and the azimuthal magnetic field converts this outward momentum into forward momentum. When the azimuthal field is strong enough, it can send the electron back into the laser

beam, where it will gain more energy from the ponderomotive force and be pushed out of the beam once again. Pukhov and Meyer-ter-Vehn [119] showed that for a high power case, $P = 6P_c$, only the head of the laser pulse is modulated with the plasma wavelength. They also observed that only a few plasma wave periods exist, indicating SMLWFA. This region contains electrons with energies of up to 30 MeV. An unmodulated region in the phase space, corresponding to the tightly focused tail of the laser pulse, where no regular wake exists, was attested to the B-loop mechanism, with an azimuthal magnetic field reaching up to $0.15B_0$. This region contains electrons with energies of up to 15 – 20 MeV.

The B-loop mechanism dominates when the regular self-modulated wake field cannot grow, this happens at densities near n_c , and when the wake field disappears, for example due to a strong wave-breaking.

- Stochastic acceleration

In this acceleration mechanism significant net energy gain by electrons may occur if the system is perturbed by a stochastic force (collision, turbulence) acting in the laser polarisation direction. One way, presented by Tanimoto *et al.* [120], is based on phase-jump disturbances of the driving laser fields. Another method of introducing stochastic laser fields has been studied by introducing some artificial anisotropic friction terms into the equations of motion [24], and is described below. The basic idea behind both methods is to introduce (small) perturbations to a plane wave, which result in a significant net energy transfer to the electrons. This occurs by dephasing the electron oscillations with respect to the laser.

From equations (A.18) one can obtain the two constants of motion:

$$p_x - a = p_{x0} \quad (2.30)$$

$$\gamma - p_z = \gamma_0 - p_{z0} \quad (2.31)$$

where $a = a_0 \sin(\omega\tau)$, and p_{x0} and p_{z0} are the initial momenta of the electron. From equations (2.30) and (2.31) one can find the momenta in terms of $a(\tau)$,

where $\tau = t - z/c$. Equation (2.30) becomes:

$$p_x(\tau) = p_{x0} + a(\tau) \quad (2.32)$$

and from squaring equation (2.31), substituting $\gamma^2 = 1 + p_x^2 + p_z^2$, solving for γ and putting the result back into equation (2.31), one obtains:

$$p_z(\tau) = \frac{1 + (a(\tau) + p_{x0})^2 - (\gamma_0 - p_{z0})^2}{2(\gamma_0 - p_{z0})} \quad (2.33)$$

From the above solutions, one notices from equations (2.30) and (2.31) that as soon as the pulse has gone ($a(\tau) = 0$), the electron momenta fall back to their initial values, $p_x = p_{x0}$ and $p_z = p_{z0}$. There is no net energy transfer, even though the electron has shifted (Lawson–Woodward criterion [118]).

Slightly perturbing the solution given above, following [24], by adding a friction term, gives the following equation of motions:

$$\frac{dp_x}{dt} = \frac{da}{dt} - \nu_{\perp} p_x \quad (2.34)$$

$$\frac{dp_z}{dt} = -v_x \frac{da}{dz} - \nu_{\parallel} p_z \quad (2.35)$$

where $\nu_{\parallel}$ and $\nu_{\perp}$ are the friction constants in the z - and x -direction, respectively. The transverse momentum is then:

$$p_x \approx a_0 \sin(\omega\tau) + \frac{\nu_{\perp}\gamma}{\omega} a_0 \cos(\omega\tau) \quad (2.36)$$

and the longitudinal equation of motion becomes:

$$\begin{aligned} \frac{dp_z}{dt} \approx & \frac{\omega a_0^2}{2\gamma} \sin(2\omega\tau) \\ & + \nu_{\perp} a_0^2 \cos^2(\omega\tau) - \nu_{\parallel} p_z \end{aligned} \quad (2.37)$$

Averaging over the laser cycle gives:

$$\left\langle \frac{dp_z}{dt} \right\rangle \approx \frac{\nu_{\perp} a_0^2}{2} - \nu_{\parallel} \langle p_z \rangle \quad (2.38)$$

This shows a force acting in the x -direction leading to a rate of change of momentum in the z -direction. While friction in the z -direction just leads to the usual damping.

Only the forces acting perpendicular to the laser propagation trigger accelerations in the laser direction. The maximum energy that the electrons can gain exceeds the ponderomotive energy ($E_p = mc^2 a_0^2/4$). Meyer-ter-Vehn and Sheng, [24], found that the highest electron energies as well as the effective temperature scale roughly as $\propto a_0$ and therefore $\propto \sqrt{I_0}$. The resulting electron spectra shows a two-temperature distribution. The lower energetic part of the spectrum has a lower temperature and is probably due to the stochastic force. Only the electrons with energies larger than 5 MeV display the above mentioned intensity scaling. The stochastic force plays a catalytic role in liberating the quiver energy intrinsic to the electrons in the laser field, and therefore leading to the hot electron tail in the spectra.

The perturbations could happen in the form of collisions or fluctuating electromagnetic fields.

2.3 Laser interactions with overdense plasma

In this section the absorption mechanisms at different intensities are described. In the regime of moderate intensities, $10^{12} \text{ Wcm}^{-2}\mu\text{m}^2 < I\lambda_\mu^2 < 10^{17} \text{ Wcm}^{-2}\mu\text{m}^2$, the classical absorption mechanisms of inverse Bremsstrahlung and resonance absorption are dominant. For higher intensities, $I\lambda_\mu^2 > 10^{18} \text{ Wcm}^{-2}\mu\text{m}^2$, vacuum heating and $\mathbf{j} \times \mathbf{B}$ heating become the dominant processes, as the electron motion becomes relativistic. The description of the absorption mechanisms is followed by a discussion of the electron transport in a laser produced plasma.

2.3.1 Absorption mechanisms

At laser intensities of $I \approx 10^{12} - 10^{14} \text{ Wcm}^{-2}$ collisions dominate the absorption physics. The dominant process depends on the density profile. Pulses with nanosecond duration are absorbed mainly by inverse Bremsstrahlung in the underdense, coronal plasma, extending out many hundreds of μm from the target surface ($L/\lambda \sim 100$, here $L = c_s\tau$ is the scale length of the density profile, where $c_s = \sqrt{k_B(ZT_e + T_i)/m_i}$ is the ion sound speed, and τ is the duration of the ionising laser.). On the sub-picosecond timescale, where $L/\lambda \leq 0.1$ is expected, a steep density profile gives access to much higher densities, where the plasma is highly collisional [78].

Collisional absorption is reduced for $I\lambda_\mu^2 > 10^{15} \text{ Wcm}^{-2}\mu\text{m}^2$, as the plasma temperature rises sufficiently fast that collisions become ineffective during the interaction ($T_e \propto I_a^{4/9}t^{2/9}$, where I_a is the absorbed laser intensity in 10^{15} Wcm^{-2} , and t is the time in the interaction in 100 fs.) [121]. At these intensities the electron quiver velocity becomes comparable to the thermal velocity, thus reducing the effective collision frequency further [122, 123].

2.3.1.1 Inverse Bremsstrahlung

Inverse Bremsstrahlung, also known as collisional heating, is the dominant absorption mechanism for $I\lambda_\mu^2 \lesssim 10^{15} \text{ Wcm}^{-2}\mu\text{m}^2$. As an electron oscillates in the electric field of the laser, energy can be transferred from the laser to the plasma, when this

electron collides with an ion [124].

Following the derivation in appendix A.2 – but adding a damping term that is proportional to the electron-ion collision frequency, $\nu_{ei} \propto n_e Z_i T_e^{-3/2}$ [125], to equation (A.28)

$$\rightarrow m_e \frac{\partial \mathbf{v}}{\partial t} + m_e \nu_{ei} \mathbf{v} = -e \mathbf{E}$$

– gives the dispersion relation for inverse Bremsstrahlung:

$$c^2 k^2 = \omega_0^2 - \frac{\omega_0}{\omega_0 + i \nu_{ei}} \omega_p^2 \quad (2.39)$$

In the coronal plasma, where inverse Bremsstrahlung takes place, $\nu_{ei} \ll \omega_0$. Therefore the dispersion relation can be approximated, using the expansion $(1 \pm x)^{-1} \approx 1 \mp x + x^2 \mp x^3$, to:

$$\frac{c^2 k^2}{\omega_0^2} \approx 1 - \frac{\omega_p^2}{\omega_0^2} + i \nu_{ei} \frac{\omega_p^2}{\omega_0^3}$$

Using the definition of the refractive index from equation (2.13), $\eta_R = \frac{kc}{\omega_0} = \sqrt{1 - \frac{n_e}{n_c}}$, and taking only the imaginary part of the equation: $k = \frac{\omega_0^2}{kc^2} - \frac{\omega_p^2}{kc^2} + \frac{i \nu_{ei} \omega_0 \omega_p^2}{kc^2 \omega_0^2}$, yields the spatial damping rate of the light wave, or the collisional absorption coefficient, κ_{IB} :

$$\kappa_{IB} = \frac{\nu_{ei}}{c} \frac{n_e}{n_c} \left(1 - \frac{n_e}{n_c} \right)^{-1/2} \quad (2.40)$$

$$\kappa_{IB} \propto \frac{Z_i n_e^2}{T_e^{3/2}} \left(1 - \frac{n_e}{n_c} \right)^{-1/2} \quad (2.41)$$

This shows that inverse Bremsstrahlung is largest for low electron temperatures, high densities and high Z_i plasmas [81, 125]. The laser can only propagate in plasma densities up to the critical density, n_c . Around this critical surface most of the inverse Bremsstrahlung absorption happens. The absorption through inverse Bremsstrahlung is better for shorter laser wavelength, higher density and lower electron temperature at the critical surface.

As the intensity of the laser increases the electron quiver velocity becomes larger which leads to a drop in ν_{ei} . Therefore κ_{IB} will drop off as the intensity increases and inverse Bremsstrahlung absorption will become less important.

2.3.1.2 Resonance absorption

Resonance absorption is a collisionless mechanism in which laser energy is transferred to plasma waves. When a p-polarised electromagnetic wave is obliquely incident on an increasing plasma gradient, which extends to above the critical density (n_c) then the light is specularly reflected at a density that is lower than n_c . The electric field component parallel to the density gradient can “tunnel through” to the critical density (where $\omega_p = \omega_0$) and excite resonant plasma waves in this region. The process of resonance absorption is schematically shown in figure 2.7.

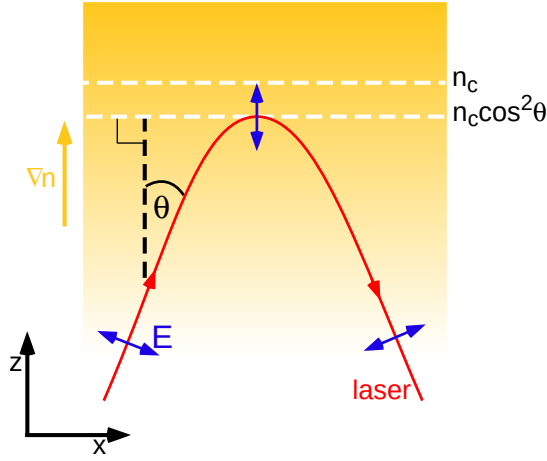


Figure 2.7: Resonance absorption mechanism for p-polarised light, obliquely incident to the plasma gradient at an angle θ .

The density at which the wave is reflected can be derived by looking at its propagation in a density ramp. Therefore one considers the dispersion relation. Since the light is incident with an angle θ , in a density ramp in the z -direction, the dispersion relation can be written as:

$$\omega_0^2 = \omega_p^2(z) + c^2 (k_z + k_x)^2$$

$$\omega_0^2 = \omega_p^2(z) + c^2 (k_z^2 + k_0^2 \sin^2 \theta)$$

Using $\omega_0 = ck_0$ and noting that the propagation stops when $k_z = 0$,

$$\omega_p^2 = c^2 k_0^2 (1 - \sin^2 \theta) = c^2 k_0^2 \cos^2 \theta$$

Therefore the light is reflected when:

$$\frac{\omega_p^2}{\omega_0^2} = \frac{n_e}{n_c} = \cos^2 \theta$$

This means that a small angle of incidence will propagate further. However it will have a smaller electric field component normal to the target, which reduces absorption. The excited plasma wave grows over a number of laser periods and is eventually damped by either collisions (at low intensities), or by particle trapping and wave-breaking (at high intensities).

The fraction of the energy transferred to the plasma wave by resonance absorption, f_{ra} is [125, 126]:

$$f_{ra} \approx \frac{\Phi^2(\tau)}{2}$$

$$\text{with } \Phi(\tau) \approx 2.3 \tau \exp\left(-\frac{2\tau^3}{3}\right) \text{ and } \tau = \left(\frac{\omega_0 L}{c}\right)^{1/3} \sin \theta,$$

where L is the density scale length $L \simeq c_s \tau_{Laser}$. The electrons accelerated by this mechanism feature a roughly maxwellian tail, superimposed on the initial background temperature. Via simulations the temperature of the hot tail was found to scale as [127]:

$$T_{hot}^{ra} \approx 14 \left(I_{16} \lambda_\mu^2 T_{keV}\right)^{1/3} \text{ [keV]} \quad (2.42)$$

where I_{16} is the laser intensity in 10^{16} Wcm^{-2} , λ_μ is the laser wavelength in μm and T_{keV} is the background electron temperature in keV.

2.3.1.3 Vacuum heating or Brunel heating

This absorption mechanism was first discussed by Brunel in 1987 [128]. It is a non-resonant absorption heating mechanism for an intense electromagnetic wave incident obliquely onto an overdense plasma with a sharp boundary. In the oscillating laser pulse, electrons can be dragged out into the vacuum and then sent back into the plasma as the fields change direction. The returning velocity of these electrons is approximately the quiver velocity and they enter the overdense plasma having gained kinetic energy directly from the laser. These electrons escape the influence of the laser, when they pass the skin depth, $l_s = c/\omega_p$.

One model frequently used to describe this mechanism is the capacitor approximation, which is described in appendix A.8 [78]. The fractional absorption rate, including relativistic effects for intensities $\geq 10^{18} \text{ Wcm}^{-2}$ takes the form:

$$f'_{bh} = \frac{I'_a}{I_x} = \frac{b^3}{2\pi} a_0 \frac{\sin^3 \theta}{\cos \theta} = \frac{b}{\pi a_0} \left[\sqrt{1 + b^2 a_0^2 \sin^2 \theta} - 1 \right] \frac{\sin \theta}{\cos \theta} \quad (2.43)$$

where $b = [1 + \sqrt{1 - f_{bh}}]$, and $f_{bh} = \frac{4}{\pi} a_0 \frac{\sin^3 \theta}{\cos \theta}$, and θ is the angle of incidence.

Further assuming that $a_0 \gg 1$, then $ba_0 \sin \theta \gg 1$ in this strongly relativistic case, and the ratio of absorbed power to incident power becomes:

$$f_{bh}^{rel} = \frac{4\pi\alpha''}{(\pi + \alpha'')^2} \quad (2.44)$$

with $\alpha'' = \frac{\sin^2 \theta}{\cos \theta}$.

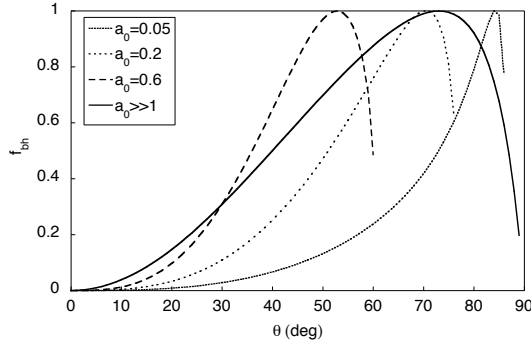


Figure 2.8: The angular dependence of the absorption rate for Brunel heating.

Figure 2.8 shows the angular dependence of Brunel heating for the non-relativistic (see equation (2.43)), as well as for the relativistic case (see equation (2.44)). The maximum of the fractional absorption rate for $a_0 < 1$ is 33%.

2.3.1.4 $\mathbf{j} \times \mathbf{B}$ heating

Unlike resonance or vacuum heating, the $\mathbf{j} \times \mathbf{B}$ heating occurs most efficiently at normal incidence to the plasma. For the $\mathbf{j} \times \mathbf{B}$ heating mechanism the intensity of the laser light has to be high enough to make the electrons move relativistically in the fields, i.e. the electrons have $v_{os} \approx c$. The magnetic field component of the

Lorentz force becomes comparable to the electric field component. At the vacuum plasma boundary, the electrons oscillating in the electric field will be directed into the forward direction by the magnetic field twice during a single laser oscillation. This leads to bunching of the electrons at $2\omega_0$ from $\mathbf{j} \times \mathbf{B}$ heating [129].

Consider a linearly polarised wave propagating in z , $\mathbf{E} = E_0(z)\hat{\mathbf{x}}\sin(\omega t)$, with the ponderomotive force, $\mathbf{F}_P = -\frac{q^2}{2\langle\gamma\rangle m_e \omega_0^2} \nabla \langle E^2 \rangle$. Then the linearly polarised wave gives rise to a longitudinal force term:

$$f_z = -\frac{m_e c^2}{2\langle\gamma\rangle} \frac{\partial \langle a^2 \rangle}{\partial z} (1 - \cos 2\omega t) \quad (2.45)$$

where a_0 is the normalised vector potential. The first term on the right hand side is the usual DC ponderomotive force, which in this context tends to push the electron density profile inwards (hole boring). The high frequency component leads to heating for any polarisation apart from circular, for which it disappears. For circular polarisation, the longitudinal force term is only the first term on the right hand side of the equation above, representing a forward push of the electrons.

The effective temperature of the hot electrons, as estimated by Wilks [53, 54] is:

$$\begin{aligned} T_{hot} &= mc^2 (\langle\gamma\rangle - 1) = mc^2 \left[\sqrt{1 + a_0^2} - 1 \right] \\ &= 0.511 \left[\sqrt{1 + 0.74 I_{18} \lambda_\mu^2} - 1 \right] \text{ MeV} \end{aligned} \quad (2.46)$$

2.3.2 Hole boring

The primary cause for the process of hole boring is a pressure imbalance – $P_L \gg P_e$, where P_L is the light pressure and P_e is the plasma pressure – causing the plasma to push inwards [54]. The ponderomotive force first accelerates and pushes the plasma electrons inwards as the laser pulse propagates into the plasma, the resultant space-charge field then also drags along the ions. As a result, the plasma in front of the laser pulse is set into motion.

Moving with the reflection front, and assuming total light and plasma reflection from opposite sides of the front, following [130] with incoming and reflected plasma density, n_i and n_r , and velocity, u_i and u_r , then number and momentum conservation

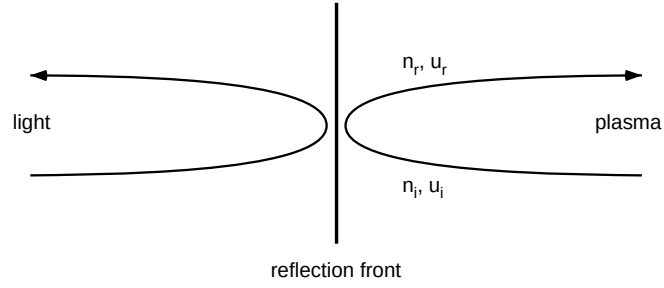


Figure 2.9: Sketch of a simple model of the light–plasma interaction, in a frame moving with the reflection front, as described in [130].

give:

$$n_i u_i = n_r u_r$$

$$m_i (n_i u_i^2 + n_r u_r^2) = \epsilon_0 E_L^2 \quad (2.47)$$

where $\frac{2I_0}{c} = \epsilon_0 E_L^2$ is the light pressure. Assuming that $u_r \sim u_i$ the velocity of the “hole boring” front can be estimated using equation (2.47), which is a balance of the momentum flux of the mass flow with the light pressure:

$$\frac{u}{c} = \sqrt{\frac{(2 - \nu_a) \cos \theta}{4} \frac{n_c}{n_e} \frac{Z m_e}{m_i} a_0^2} \quad (2.48)$$

Here Z and m_i are the charge and mass of the ions. Equation (2.48) is similar to the one found by Wilks *et al.* [54], but was modified to include the effect of absorption ν_a , and an angle of incidence θ by Zepf *et al.* [131].

Experimentally the velocity of the ion front can be measured by looking at the frequency shift of the reflected laser.

$$\omega_{shift} = \omega_0 \frac{1 - (u/c)}{1 + (u/c)}$$

Zepf *et al.* [131] studied the intensity dependence of u , and found excellent agreement with equation (2.48), assuming 50% absorption.

2.3.3 Self induced transparency

An increase in the relativistic inertia of the electrons for $a_0 \gg 1$, leads to a correction to ω_p (because ω_p depends on m_e). Therefore the critical density n_c increases to: $n_{c\gamma} = \langle \gamma \rangle n_c$, where for linearly polarised light $\langle \gamma \rangle = \sqrt{1 + \frac{a_0^2}{2}}$. Resulting in light being able to propagate to densities above the initial n_c . This induced relativistic optical transparency means that plasma with densities $n_e < n_{c\gamma}$ is relativistically transparent [84, 132–134].

This self induced transparency was found to influence the hole boring distance, d , (for $n_e < 10n_c = n_{c\gamma}$) in the following way [135]: Assuming that the laser energy –

$$\mathcal{E}_L = \tau_L A I_0 = \tau_L A \frac{c\epsilon_0}{2} E_L^2 = \tau_L A c \epsilon_0 \left(\frac{m_e c \omega_0}{e} \right)^2 \langle a^2 \rangle,$$

with A the area of the focal spot – is mostly converted into hot electrons, and a constant absorption rate – where the average energy lost to each electron is the ponderomotive energy,

$$(\langle \gamma \rangle - 1) m_e c^2 = \frac{1}{2} \langle a^2 \rangle m_e c^2$$

– the total energy lost to a channel of length d will then be:

$$\mathcal{E}_p = \frac{1}{2} \langle a^2 \rangle m_e c^2 n_e A d.$$

Equating the two energies and solving for d gives

$$d = 2c\tau_L n_c / n_e.$$

This shows that when the electron density is decreased, the laser is able to penetrate further, making a longer cavity in the plasma.

2.3.4 Fast electron transport

Alfvén [136] noted that the main limiting factor on the propagation of an electron beam in a conductor is the self-generated magnetic field, which acts to turn the electrons back towards the source, limiting the current to a value of the order of [137]:

$$\begin{aligned} I_A &= \frac{4\pi}{e\mu_0} p = \frac{1}{e\mu_0\epsilon_0} \gamma \beta m_e c = \frac{m_e c^3}{e} \beta \gamma \\ &= 17\beta\gamma \quad [\text{kA}] \end{aligned} \tag{2.49}$$

where p is the electron momentum, and $\beta = v/c$. Above this Alfvén limit, the azimuthal magnetic field around the electron beam increases so much that the Larmor radius, $r_L = \gamma m_e v_\perp / qB$, becomes smaller than the beam radius ($r_L < r_B$), and the electrons are no longer able to propagate.

However, due to the colder background electrons in the target, trying to restore current neutrality, a return current, $\mathbf{j}_{return}$, is set up, which also leads to a reduction of the magnetic field resulting from the fast electron beam. Therefore the fast electron current can be greater than the Alfvén limit ($I_{max} > I_A$). The total current is given by: $\mathbf{j}_{total} = \mathbf{j}_{fast} + \mathbf{j}_{return}$. Because the response of the return current is very fast, one can neglect the displacement current term ($\epsilon_0 \mu_0 \frac{\partial \mathbf{E}}{\partial t}$) in Ampère's law and obtains:

$$\nabla \times \mathbf{B} = \mu_0 (\mathbf{j}_{fast} + \mathbf{j}_{return})$$

Substituting $\mathbf{E} = \eta \mathbf{j}_{return}$ gives the electric field induced by the fast electrons:

$$\mathbf{E} = \frac{\eta}{\mu_0} \nabla \times \mathbf{B} - \eta \mathbf{j}_{fast}$$

where η is the resistivity [138]. The magnetic field growth is:

$$\frac{\partial \mathbf{B}}{\partial t} = -\nabla \times \mathbf{E} = \nabla \times \eta \mathbf{j}_{fast} - \nabla \times \left(\frac{\eta}{\mu_0} \nabla \times \mathbf{B} \right)$$

The first term on the right hand side corresponds to the source term, and the second term is the resistive diffusion of the magnetic field term.

It was shown, by Bell and Kingham [139], that fast electrons can be collimated into a beam, and the collimation occurs, when $R/r_L > \theta_{1/2}^2$, where R is the radius of the beam, r_L is the fast electron Larmor radius and $\theta_{1/2}$ is the half angle of the divergence of the electron beam. In this case, the magnetic field is sufficient to bend the fast electron trajectory through an angle $\theta_{1/2}$ in the distance $R/\theta_{1/2}$ in which the beam radius approximately doubles. They found that the collimation is governed by resistivity. As the target is heated the resistivity decreases and the collimation is self-limiting. Large beam power leads to weak collimation, due to the weak resistivity as the cold plasma is rapidly heated and the magnitude of the electric field, needed to draw a return current, is reduced.

2.3.5 Recirculation

Hot electrons are generated at the front of the target. As the first electrons leave the target, a sheath field is set up, and the following electrons can be reflected by this field. They return into the target, and, depending on the target thickness, undergo a number of recirculations in the target. The basic idea is illustrated in figure 2.10.

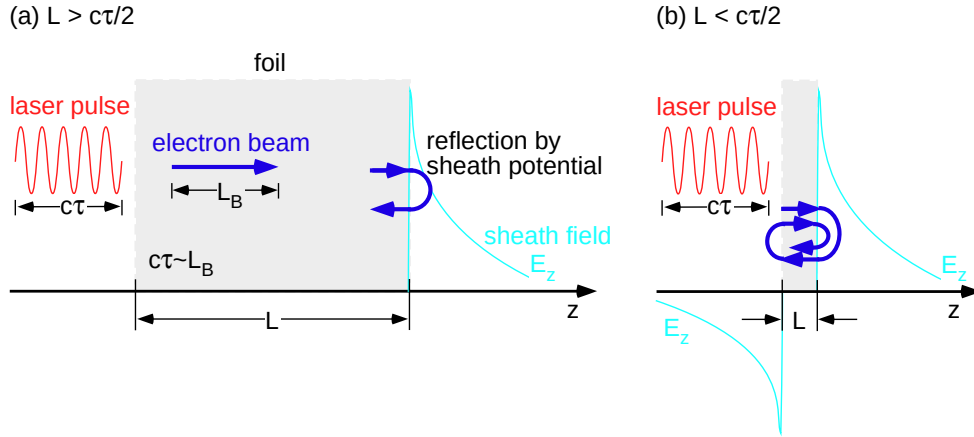


Figure 2.10: Model of recirculation for (a) thick targets, where $L > c\tau/2$, and for (b) thin targets, where $L < c\tau/2$.

The hot electron bunch has approximately the same length as the laser pulse ($L_B \sim c\tau$), since the velocity of the electrons is approximately the speed of light. They travel through the solid target without energy loss, and are then reflected at the target surface by a self-induced sheath field. When the target is thick, i.e. the target thickness L is larger than half the laser pulse length, $L > c\tau/2$, as illustrated in figure 2.10(a), the electrons overlap only locally at the target edge, and no significant increase of the hot electron density is possible. However, if the target is thin, i.e. $L < c\tau/2$ (figure 2.10(b)) the recirculation of the electrons increases the hot electron density within the target.

PIC simulations by Makinon and Sentoku [69, 70] showed that recirculation increases the number of hot electrons in the target, but does not increase the temperature [69, 70, 140].

Following the derivation of the sheath electric field in [141] (see appendix A.9),

the following scaling can be obtained:

$$E_z(\xi) = \frac{k_B T_e}{e \lambda_{D\text{hot}}} \left(1 - \frac{\xi}{\sqrt{1 + \xi^2}} \right)$$

where $\lambda_{D\text{hot}} = \sqrt{\frac{2\epsilon_0 k_B T_e}{n_{Q0} e^2}}$, $\xi = z/r_B$ and $r_B = r_L + d \tan \theta$. Here r_B is the radius of the electron beam at the rear of the target, r_L is the radius of the laser spot, d is the target thickness and θ is the half-angle of the electrons traveling through the target.

Recirculation leads to an enhancement of the proton acceleration from the rear target surface by the electrostatic sheath set up by the hot electrons. This has been experimentally observed in [70], where both the proton temperature and the cutoff energy increase with decreasing target thickness, because of an increased hot electron density, caused by recirculation.

Chapter 3

Experimental Methods

This chapter describes the experimental methods used to obtain the results reported in this thesis. The principal diagnostic for all the experiments is an electron spectrometer used to detect the accelerated electrons. As well as the energy spectrum, measurements were also made of the electron beam profile. This chapter also discusses the characterisation of the target used in these experiments, namely a supersonic gas jet, and includes a brief overview of the VULCAN laser system.

3.1 Electron diagnostics

This section describes the electron spectrometer that was used in the experiments described in chapters 4 to 6, and explains how both the energy and number of electrons detected are calibrated.

3.1.1 Spectrometer

In the experiments described, the electron spectrometer is usually situated along the laser axis. Figure 3.2 shows a schematic and figure 3.3 shows an image of the electron spectrometer used most often in the experiments described in this thesis. As accelerated electrons pass through a magnetic field of the spectrometer, the Lorentz force deflects the electrons away from their axis of propagation.

$$\frac{d\mathbf{p}}{dt} = \frac{q}{\gamma m_e} \mathbf{p} \times \mathbf{B}$$

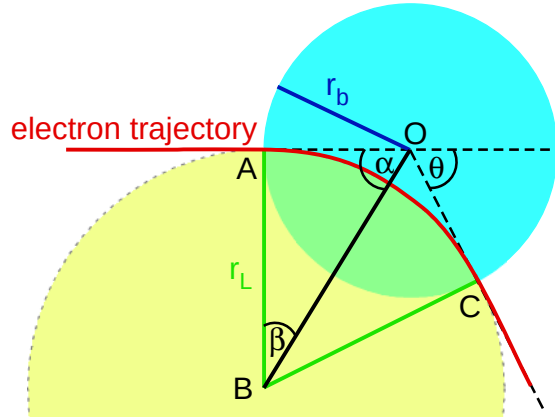


Figure 3.1: Trajectory of an electron passing through an ideal (top-hat profile) on-axis magnetic spectrometer. The light blue circle represents the magnetic field region, radius r_b . The circumference of the larger yellow circle represents the Larmor orbit of the electron, radius r_L .

In the following, we assume the electron is travelling in the z -direction and the magnetic field is in the y -direction. For a uniform magnetic field within a circular region of radius r_b with zero field outside this region given by:

$$\mathbf{B}_y(r) = B_{max} \mathbf{e}_y \quad r \leq r_b \quad (3.1)$$

$$= 0 \quad r > r_b \quad (3.2)$$

there is an analytical solution for the angular deflection of electrons passing through this field. Figure 3.1 indicates the electron trajectory and the magnetic field position used in this calculation. The path of a relativistic electron within the magnetic field is the Larmor orbit in the z - x plane having a radius given by:

$$r_L = |\mathbf{p}|/eB_{max} \quad (3.3)$$

By considering similar triangles (OBA and OBC), we can see that the angle of deflection, θ , is related to the angle α by:

$$2\alpha + \theta = \pi \quad (3.4)$$

Since the angle α is related to the magnetic field radius and the Larmor radius through:

$$\tan \alpha = \frac{r_L}{r_b} \quad (3.5)$$

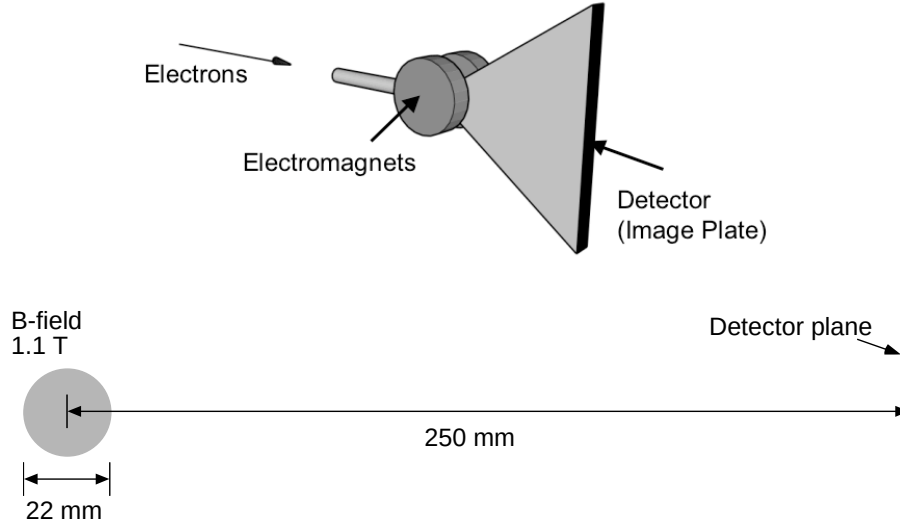


Figure 3.2: Schematic of electron spectrometer and an illustration of the parameters chosen for the comparison of the tracking codes with the analytical solution.

it follows that the angular deflection θ of an electron with momentum $|\mathbf{p}|$ is given by:

$$\tan\left(\frac{\theta}{2}\right) = \frac{r_b}{r_L} = \frac{eB_{max}r_b}{|\mathbf{p}|} \quad (3.6)$$

where r_b is the radius of the magnetic field region and r_L is the Larmor radius of the electron.

The magnetic field is induced by a current running through coils around the two circular pole pieces. These pole pieces are separated by a variable distance. A yoke surrounding the coils prevents unwanted fringe fields which would reduce the dispersion of the magnet due to regions of field with reversed polarity.

In order to convert the electron spectra from the image plate data, a tracking code written in Matlab by the author was used. Before it was used for data analysis it was tested and compared to the analytical solution, as well as an old tracking code, written in IDL. The IDL tracking code is described in [142]. The parameters chosen for the comparison with both the analytical solution and the IDL code were: A field of radius 11.0 mm with a magnetic field of 1.1 T. The deflection of the electrons was recorded at a distance of 250 mm from the centre of the circular magnetic field

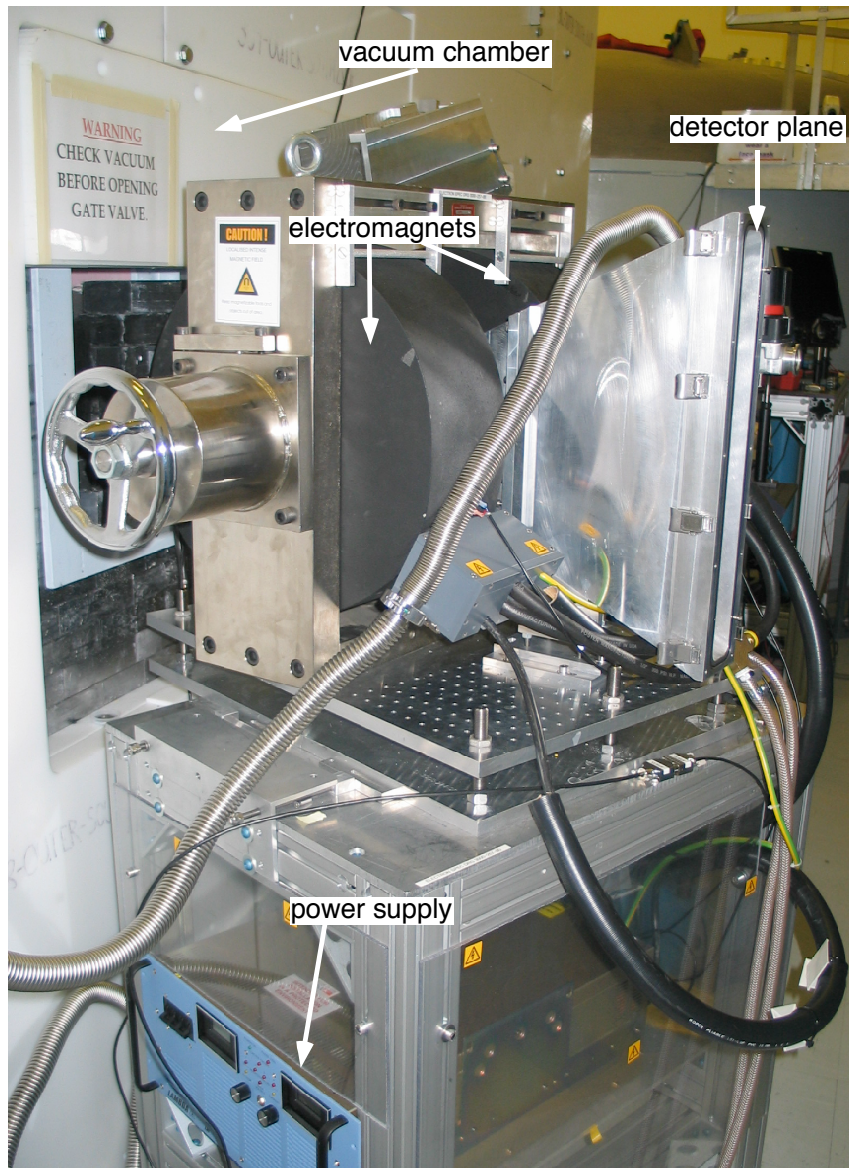


Figure 3.3: Picture of the electron spectrometer used in most of the experiments discussed.

region. Figure 3.4 shows four plots; 3.4(a) shows the very good agreement of the Matlab code and the analytical solution. In figure 3.4(b) the difference between analytical and tracking codes for both, the Matlab tracking code, as well as the IDL tracking code, is plotted. One can easily see that the accuracy has improved drastically.

Just like the IDL code, the Matlab tracking code tracks a set of different energy electrons through the magnetic field. However, in the Matlab tracking code the electrons are tracked until they reach the detector plane. This is done by solving the Lorentz force numerically using a centre-differencing method. The algorithm calculates the electron momenta at time t from the momenta and positions at the two previous time steps $t_{-1} = t - dt$ and $t_{-2} = t - 2dt$.

$$\begin{aligned}
 p_z &= p_{z-2} - e \, 2dt \, v_{x-1} \, B(z-1, x-1) \\
 p_x &= p_{x-2} + e \, 2dt \, v_{z-1} \, B(z-1, x-1) \\
 v_{z,x} &= p_{z,x} / (m_e \gamma) \\
 z &= z_{-1} + v_{z-1} * dt \\
 x &= x_{-1} + v_{x-1} * dt
 \end{aligned} \tag{3.7}$$

with the electrons propagating in z -direction, and $\gamma = \frac{KE(MeV)}{m_e(MeV/c^2)c^2} + 1$.

The code also includes the effect of beam divergence. This is done by tracking three electron trajectories for each energy, one on axis, and the other two at the angles $+d\phi$ and $-d\phi$. In the experiments the angle $d\phi$ is set by the collimator used, and is the limiting factor for the energy resolution of the spectrometer. As these three electron trajectories experience different path lengths within the magnetic field, the diverging beam can be focused by the magnet. This focusing occurs at a different position for each energy. The tracking code only considers the 2-dimensional profile of the field as this is sufficient to determine the energy dispersion curves. Both the analytical code, and the test version for the tracking code can be found in appendix B.

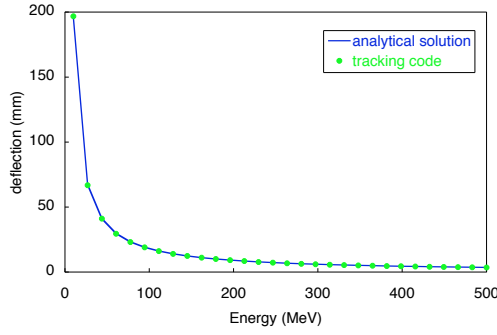
Figure 3.4(c) shows the three dispersion curves for a typical set-up of the electron spectrometer as used in the experiment described in section 6.2. The deflection-

distance between source and magnet centre	2492 mm
distance between magnet centre and detector plane	500 mm
size of magnet pole pieces	152.4 mm
magnetic field strength	0.067T
collimator half angle	1.96 mrad

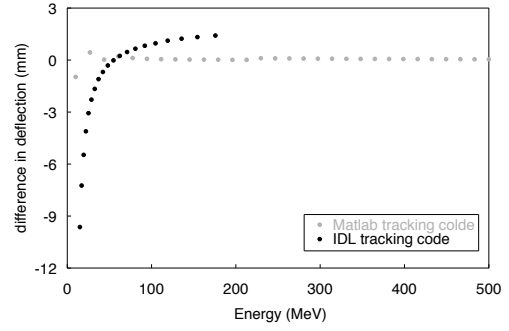
Table 3.1: Parameters used to generate the dispersion curves in figure 3.4(c).

energy relation for the magnet is of the form $E = ax^b$ where $b \simeq -1$, the dispersion dE/dx is therefore abx^{b-1} .

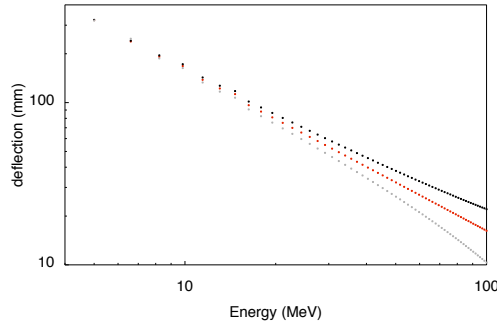
The electron spectrometer used in most experiments described in this thesis is the electron spectrometer pictured in figure 3.3. It uses an electromagnet manufactured by AML Inc. with pole pieces of diameter 152.4 mm, and 101.6 mm pole caps. The gap between the pole pieces is set to 50 mm, and the magnetic field strength shows a linear dependence on the current, as shown in figure 3.4(d).



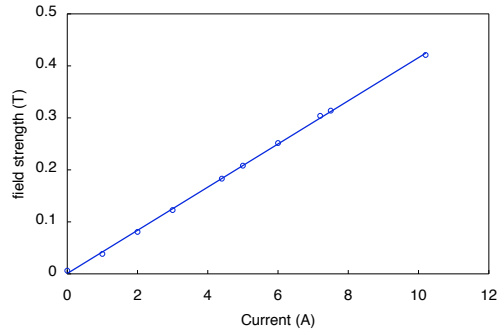
(a) Comparison of analytic solution for ideal magnetic field profile and results from the tracking code.



(b) Comparison of Matlab and IDL tracking codes. The difference between the deflection for the analytical solution and the tracking codes is plotted against the kinetic energy.



(c) Electron deflection curves for the parameters described in table 3.1. The different colours indicate different electron trajectories, on axis (red) and $\pm d\phi$ (black and grey).



(d) Magnetic field strength dependence on the current, for the range used in the experiments described in this thesis.

Figure 3.4: Comparison/testing of Matlab tracking code with analytical solution and old IDL tracking code. And electron deflection curves, as well as the dependence of the magnetic field strength on the current.

3.1.2 Detectors

Imaging plate

The electron detectors used in the experiments described were imaging plates (IP) [143, 144]. Imaging plates are a re-usable radiation detectors which produce an image of where the electrons pass through the plate. The imaging plates used in the experiments described in this thesis were the Fuji BAS-MS type. A sketch of their composite structure can be seen in figure 3.5. Imaging plates utilise a process called photo-stimulated luminescence (PSL). The plates contain a $100\ \mu\text{m}$ layer of luminescent material, BaFBr:Eu^{2+} (barium fluorbromide containing a trace of bivalent europium). The europium ions are excited into a metastable state via energy deposition in this layer. This metastable state cannot decay through fast optical transitions. Therefore the information of energy deposition is stored as a latent image in the plate. This image can be read by deliberately exciting the europium ions further, into a state that is not metastable. As the ions relax from this state they emit a photon with $\lambda_{PSL} \simeq 390\ \text{nm}$, which is measurable. This process is illustrated in figure 3.6. The number of europium ions excited is directly related to the amount of energy deposited in the plate. The reading is performed with a commercially available Fuji BAS 1800II image plate reader. The exposed image plate is scanned with a focused laser beam ($\lambda_L \simeq 633\ \text{nm}$). The photo-stimulated luminescence released by the laser is collected into a photomultiplier tube and is converted to electric signals. Because the laser scans across the plate the resolution of the imaging plate is effectively determined by the size of the focal spot of the laser and can be set to $50\ \mu\text{m}$, $100\ \mu\text{m}$ and $200\ \mu\text{m}$. In normal operation the resolution is set to $200\ \mu\text{m}$, as the resolution of the electron spectrometer is

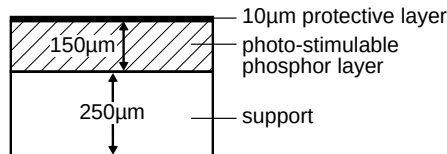


Figure 3.5: Composite structure of imaging plate (IP).

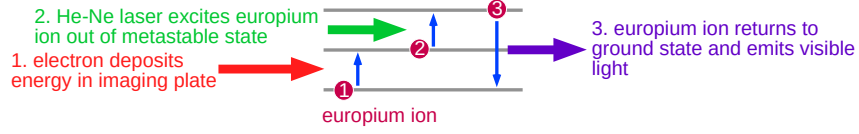


Figure 3.6: Sketch of the Photo-Stimulated Luminescence process (PSL).

limited by the collimating aperture, so no extra information is gained using a higher resolution scan. For some shots the resolution on the scanner was set to $50 \mu\text{m}$, the reason behind this is that we wanted to gain x-ray signal resolution. In the signal conversion, this is taken into account, using equation (3.8).

The scanner saves two files, *filename.img* and *filename.inf*. The BAS system from Fuji generates the image data as a logarithmic compression. Therefore it is necessary to make a log (quantum level QL) to linear (photo-stimulated luminescence PSL) conversion [145, 146]:

$$C_{PSL} = \left(\frac{Res}{100} \right)^2 \times \frac{4000}{S} \times 10^{L \times (\frac{C_{QL}}{G} - \frac{1}{2})} \quad (3.8)$$

Where C_{PSL} is the quantified value in linear scale, Res is the scan resolution in μm , S corresponds to the sensitivity setting (usually 4000), L denotes the latitude setting of the scanner (usually 5), G is the bit depth of the scan (for 8 bit=256, 10 bit=1024, 12 bit=4096 and 16 bit=65536), and C_{QL} is a pixel value in the file *filename.img*. The information on the other values (S , Res , G and L) is stored in the *filename.inf* file. Because the C_{QL} data is proportional to the logarithm of the actual C_{PSL} level it is vital that the C_{QL} is converted to C_{PSL} before any analysis such as background subtraction or noise filtering is performed.

The imaging plate response for electrons above 1 MeV is very constant, and the conversion of C_{PSL} to number of electrons for the BAS MS2325 type is $1 C_{PSL} = 50 \text{ electrons}^*$, with an error of $\pm 10\%$. Using a Monte-Carlo code, which takes into account the energy loss and deposition due to collisions and radiative processes, it was found that an electron passing through the IP deposits 50 keV in a $50 \mu\text{m}$ pixel for electrons in the range of 1 to 100 MeV [147], see figure 3.7. The conversion factor

*This calibration for a scan of the imaging plate at $50 \mu\text{m}$ resolution, was found to be $0.02 C_{PSL}$ corresponding to the deposition of 50 keV.

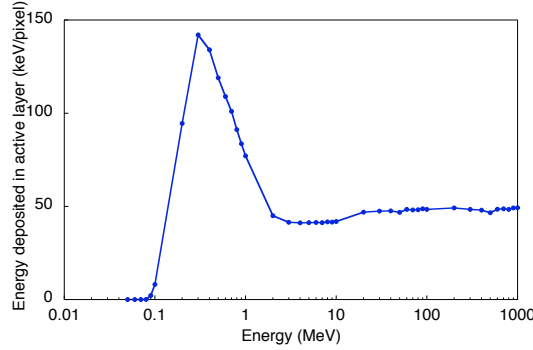


Figure 3.7: Imaging plate response for type BAS MS2325 (courtesy of AWE).

from C_{PSL} to number of electrons can change for different types of imaging plate. The value given here is only valid for the Fuji BAS-MS type, for the SR 2025 type it was found to be 1 $C_{PSL} = 140$ electrons, with an error of $\pm 10\%$ [148, 149].

3.1.3 Stacks

Using the fact that the range of an electron in matter depends on its energy, stacks can be used to discern the divergence at different energies. Electron beam divergence stacks are made out of layers of metal pieces and detectors (e.g. imaging plate). Depending on their position in the stack, the detectors are hit by electrons above a certain energy threshold. This threshold is defined by the range of the electrons in the material in front of the detector.

Extracting the desired information from these stacks is complicated. On the one hand the unfolding of the data from a matrix equation is necessary, because the signal in the detector is due to electrons with a range of energies above a certain level. Additionally it has to be considered that electrons as they pass through the stack undergo many small angle scattering events due to Coulomb scattering from the nuclei in the stack material [142, 150].

The stacks can be used to discern the divergence angle of the electron beam, and to estimate the number of electrons. No electron spectra were recovered from this kind of measurement. For a description of the unfolding process please refer to [142, 150].

3.2 Gas jets

For the experiments in laser interactions with underdense plasma, the targets were supersonic gas jets. In order to be in the underdense ($n_e < n_c$) plasma regime the target material must be gaseous. The gases used are helium and hydrogen.

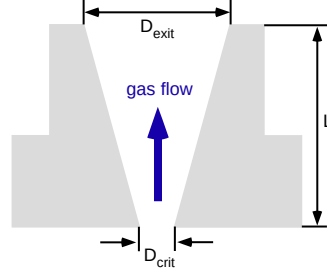


Figure 3.8: Cross section of a conical gas jet nozzle.

The reason for using a supersonic gas jet is that the supersonic flow ensures a sharp interface and a uniform density profile over a given distance, which depends on the nozzle diameter (D_{exit}) [151]. Since the gas expands after leaving the nozzle, the density decreases as the distance increases. The shape of the nozzle influences the density profile of the gas jet. Semushin and Malka, [151], found that a conical nozzle, as shown in figure 3.8, produces a flat density profile with short density ramps into vacuum. The important parameters of such a design are the aperture through which the gas enters the nozzle, D_{crit} , the nozzle diameter at the exit, D_{exit} , and the length between those two, L . The nozzles used in the experiments described in this thesis were designed using the optimised values found in [151].

Table 3.2 shows the optimum length and achievable neutral density at 100 bar backing pressure, as given in [151], as well as the lengths for the different nozzle materials as used in the experiments. However, the density of the gas jet depends on a number of factors, including the triggering time relative to the time the laser arrives. This was usually set to 15 ms. Other factors include the solenoid used, the voltage used to open it, the backing pressure of the target gas, as well as the shape of the nozzle and the nozzle diameter. The nozzles used in the experiments, which are summarised in table 3.3, were made out of brass, aluminium (run-II and

D_{exit} (mm)	D_{crit} (mm)	L_{opt} (mm)	n_{exit} (10^{19}cm^{-3})	$L_{brass/alu}$	$L_{plastic}$
10	1	15	0.75	15	-
5	1	10	2.7	10	10
3	1	7	7.5	7	7
2	1	6	18	6	5
1	0.5	4	16	4	4

Table 3.2: Parameters for the nozzles used in the experiments. The design relies on simulations conducted in [151].

$I < 10^{20} \text{ Wcm}^{-2}$) and plastic (run-I). The solenoid valves used were the Parker series 9 (1.5 mm orifice) for the experiments in the intensity regime above $1 \times 10^{20} \text{ Wcm}^{-2}$. For the experiment in the $I < 10^{20} \text{ Wcm}^{-2}$ regime a Peter Paul valve (model EH22) with an orifice of 1.2 mm was used. The electron density n_e during the data shots was measured using the frequency separation of the forward Raman scattered laser spectra, as described in section 2.2.2, using $\Delta\omega = \omega_p = \sqrt{\frac{n_e e^2}{m_e \epsilon_0}}$.

Using interferometric techniques the gas density profile can be measured. This is usually done using argon as the backing gas, and measuring the neutral density [152]. This is justified since it has the same adiabatic index as helium. A density map can be created by comparing an interferogram image with a reference image of the fringes without gas. Applying an Abel inversion, which assumes that the gas jet is cylindrical, creates a two-dimensional density map of the gas jet. The neutral density profile is found by taking a line out at the height above the gas jet at which the interaction with the laser takes place. The analysis described above can be done using the software IDEA [153].

3.3 The Vulcan laser system

The experiments described in this thesis were conducted with the VULCAN laser, which is part of the Central Laser Facility (CLF) at the Rutherford Appleton Laboratory (RAL), Oxfordshire, UK. It is a Nd:glass laser system, with a wavelength of $1.05 \mu\text{m}$ and has three distinctly different target areas, Target Area Petawatt (TAP), Target Area West (TAW) and Target Area East (TAE). Only Target Area Petawatt and Target Area West were used for the experiments discussed, and will be described here. TAP has a single beam-line capable of petawatt power ($10^{15} \text{ W} = 400 \text{ J}/400 \text{ fs}$) and can be focused down to intensities of 10^{21} Wcm^{-2} , making it one of the highest intensity laser systems in the world at present. TAW has a 100 TW, sub-picosecond beam, along with a 10 ps probe and up to 6 other beams each of 200 J in a nanosecond. For the experiment in TAW that is described in this thesis only the 100 TW and the probe beams were used and the pulse length of the main pulse was reduced to 640 fs. TAW has recently undergone an upgrade, adding a second high energy short pulse beam, delivering 500 J in 10 ps. The short pulse beam lines on VULCAN utilise the chirped pulse amplification and optical parametric amplification techniques described in the following sections.

3.3.1 Chirped pulse amplification

Optical chirped pulse amplification (CPA) was developed by Mourou and Strickland [1]. Before that, one of the major limiting factors in intensity gain was the damaging of the gain medium due to non-linear processes, in particular self-focusing in the amplifier medium. These non-linear processes grow with a rate that is determined by the B -integral: $B = \frac{2\pi}{\lambda} \int \eta_2 I(z) dz$, where $I(z)$ is the optical intensity ($I = \frac{c\epsilon_0\eta}{2}|E^2|$), z is the position in beam propagation direction and η_2 is the non-linear refractive index[†]. As $\eta_2 I$ is the non-linear change in the refractive index, the B -integral is the total on-axis non-linear phase shift accumulated in a passage through the medium. For very high intensity laser pulses the B -integral even for air would

[†]When light with high intensity propagates through a medium, this causes non-linear effects. The simplest of these is the Kerr effect, which can be described as a change (usually an increase) in the refractive index in proportion to the optical intensity I , $\Delta\eta = \eta_2 I$.

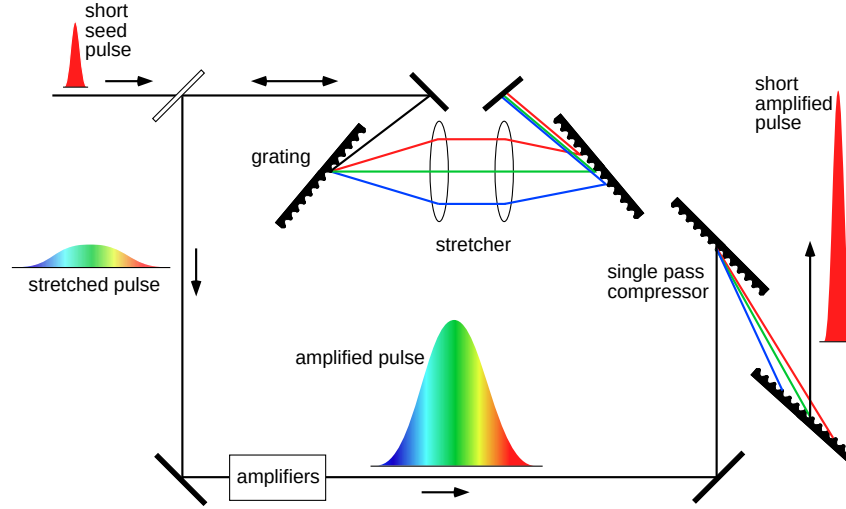


Figure 3.9: A schematic of the chirped pulse amplification (CPA) technique.

be enough to degrade the laser pulse. Therefore the interaction must take place in a vacuum chamber.

One way to avoid the non-linear effects was to reduce I by using expensive large beam diameter laser systems. Another, now widely used method, is CPA, where an ultrashort low energy laser pulse is stretched in time prior to introducing it to the gain medium, here the stretcher consists of a pair of anti-parallel gratings. This results in the low-frequency component of the laser pulse traveling a shorter path than the high-frequency component. After going through this stretcher, the laser pulse becomes positively chirped, which means that the high-frequency components lag behind the low-frequency components, and has a longer pulse duration than the original. Then the stretched pulse, which has an intensity that is sufficiently low compared with the damage threshold fluence of the amplifying medium, is safely introduced to the gain medium and amplified. Finally, the amplified laser pulse is recompressed back to the original pulse width through a reversal of the stretching process, i.e. a single pass compressor. This technique can achieve orders of magnitude higher peak power than laser systems could generate before the invention of CPA.

3.3.2 Optical parametric chirped pulse amplification

The use of CPA allowed significant improvement in the maximum intensities which could be achieved. However, 1 PW with pulse durations of several hundred femtoseconds currently represents the limits in single beam performance, resulting from the energy handling limits of the compressor gratings and the gain bandwidth of the amplifiers. The route to yet higher powers lies in further reduction of the pulse duration. This requires an amplifier with increased gain bandwidth at high energy. Which is provided by the optical parametric amplification (OPA) technique [154].

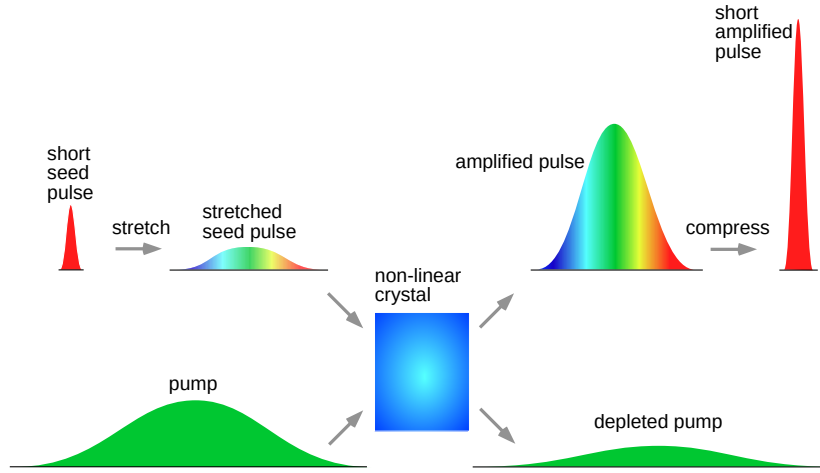


Figure 3.10: Schematic of the optical parametric chirped pulse amplification scheme.

As described in section 3.3.1, the seed pulse for the amplification is chirped. It is then mixed in a non-linear crystal, such as β -barium borate (BBO), with a pump laser. This pump laser is usually a nanosecond pulse of a few Joules. In order to fully exploit the large gain bandwidths possible in OPAs, the temporal shape of the pump pulse should be close to a top hat, with uniform intensity in space and time, as this reduces gain narrowing effects. The pulse length of the pump beam also influences the high energy extraction efficiency. In the non-linear crystal the pump energy of frequency, ω_{pump} , is parametrically converted into the seed frequency, ω_{seed} , by the process; $\omega_{pump} \rightarrow \omega_{seed} + \omega_{idler}$, where ω_{idler} is an idler beam which has to be produced for the energy conservation condition to be satisfied. Another condition that has to be met is the phase matching condition, $\mathbf{k}_{pump} \rightarrow \mathbf{k}_{seed} + \mathbf{k}_{idler}$, where

$\mathbf{k}_{pump}$ is the wave vector of the pump. The amplified pulse can be re-compressed as described in section 3.3.1.

With the OPCPA technique the amplification is constrained to a short time window, since the pump duration is approximately the same length as the signal pulse. Recently OPCPA has been performed on even shorter timescales so as to minimise the length of the parametric fluorescence and pre-pulse. The pedestal before the main pulse is non negligible also for the OPCPA amplification and is due to parametric fluorescence (which also meets the above phase matching conditions). In standard amplifiers, where the amplifying medium is a lasing material, there is a degree of amplified spontaneous emission (ASE), if the gain duration is longer than the pulse to be amplified. The contrast ratio of a laser system is the ratio between the peak intensity and the pedestal. The intensity threshold for plasma generation from solid targets lies at about 10^{10} Wcm^{-2} , and since laser systems are now capable of producing focused intensities on the order of 10^{21} Wcm^{-2} , a contrast ratio of 10^{11} or better would be required to prevent ionisation and the formation of a pre-plasma before the main pulse arrives.

3.3.3 Target Area Petawatt (TAP)

The VULCAN Petawatt was upgraded in 2002. It was specified to deliver a beam of 670 J in the stretched pulse, delivering 400 J on target in a compressed pulse of 400 fs, giving a peak irradiance of 10^{21} Wcm^{-2} in a $5 \mu\text{m}$ focal spot [2, 155–157]. To produce such a pulse, a customised stretcher and compressor, as well as a new front-end using OPCPA pre-amplification, a mixed glass amplifier (Nd:silicate and Nd:phosphate) and adaptive optics, were needed. A schematic of the system is shown in figure 3.13.

The pulse generator for the system is a commercial Ti:Sapphire oscillator operating at 1055 nm and delivering pulses of 5 nJ, 120 fs long at a repetition rate of 80 MHz. It uses the Kerr-lens mode-locking (KLM) method, and produces a bandwidth of 15 nm. Before amplification the oscillator pulse is stretched to a pulse length of 2.4 ns. Stretcher and compressor have to be coordinated in order to get

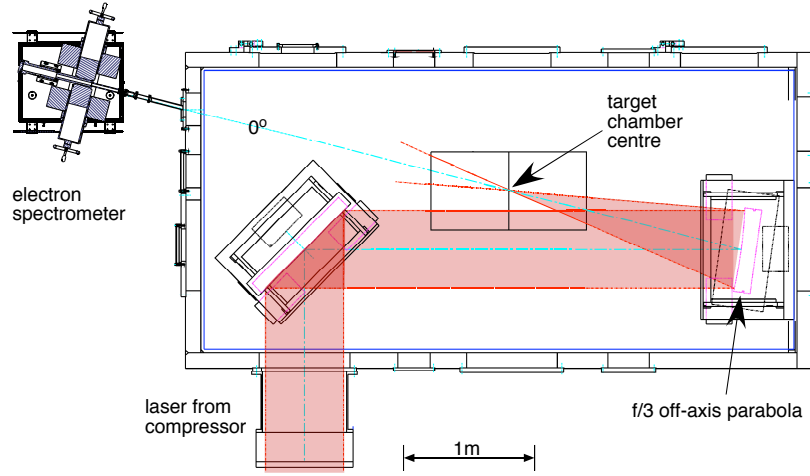


Figure 3.11: Sketch of the target chamber in TAP, including the electron spectrometer.

the shortest possible pulse.

The pre-amplifier is based on the OPCPA technique, described in section 3.3.2. The amplification of the seed is achieved in three single-pass stages. Stage one and two are designed to give a gain of 10^3 , and the third stage is used to saturate the process and gives a gain of 10^2 . BBO crystals are used as the non-linear medium for all three stages. All three stages are pumped by a Q-switched Nd:YAG. It is a Single Longitudinal Mode (SLM) laser that generates a train of 1 J pulses at 10 Hz. The pulses have a gaussian shape and a pulse-length of 15 ns at 532 nm. Pockel cells slice the central 3 ns, such that the pump pulse has an almost top hat temporal profile. The total gain from this pre-amplification stage is higher than 10^7 . After the OPCPA stage the pulse is stretched to 4.8 ns pulse length.

Phosphate glass amplifiers are the standard amplifiers used in the VULCAN system and are pumped by flash lamps. Using silicate glass as well as phosphate glass amplifiers has the advantage of broadening the bandwidth that can be amplified and hence reducing the pulse length that can be delivered after the compressor. Therefore the amplifier chain for the petawatt laser consists of a 9 mm single pass silicate rod amplifier, followed by a double pass 16 mm silicate rod, then 25 mm

and 45 mm single pass phosphate rod amplifiers followed by a double pass 108 mm phosphate disc. In addition to these amplifiers, three ex-Nova 208 mm aperture disk amplifiers are used to provide the 670 J laser energy required for the Petawatt beam line. This final amplification stage has gain isolation, by means of a 208 mm aperture Faraday rotator and a single polariser at the stage input. The diameter of the amplifiers increases, so that the fluence through each stage is below the damage threshold, but high enough so that the gain is saturated to reduce gain-narrowing of the bandwidth.

From the final amplification stage, a 19 m long vacuum spatial filter improves the beam quality and expands the beam from 208 mm to 600 mm. The gold-coated holographic gratings, of 940 mm diameter and a line density of 1480 lines/mm, were supplied by the Lawrence Livermore National Laboratory (LLNL). They are set in a single pass geometry with a separation of 13 m and an input operational angle of 48° . The compressor is housed in a vacuum chamber which can be isolated from the target vacuum chamber by a gate valve. Two mirrors are used to bring the compressed beam onto the $f/3$ off axis parabolic mirror. The first mirror turns the beam into the target chamber, the leakage through the back of this mirror is used for the beam diagnostics, which include the spectral width and auto-correlator measurements for the pulse duration. The second mirror turns the compressed beam onto the focusing optic. Figure 3.11 shows a schematic of the TAP target chamber.

An 120 mm diameter adaptive optics (AO) module [155] is used to improve the wavefront quality to obtain the optimal focal spot on target and increase the intensity. Wavefront errors arise from static aberrations of the beam-line optics and distortions caused by the thermal gradients introduced to the amplifiers by firing consecutive laser shots. Without the phase correction, VULCAN operates with a beam quality of approximately three times the diffraction limit. Improving the wavefront quality with the adaptive optic, the optimal focal spot improves by a factor of > 2 , increasing the intensity on target by a factor of 4. Wavefront error correction also helps achieve the best re-compression.

As with any high intensity laser system, pre-pulses can be present due to amplified spontaneous emission (ASE) and phase errors. For full disc shots the best

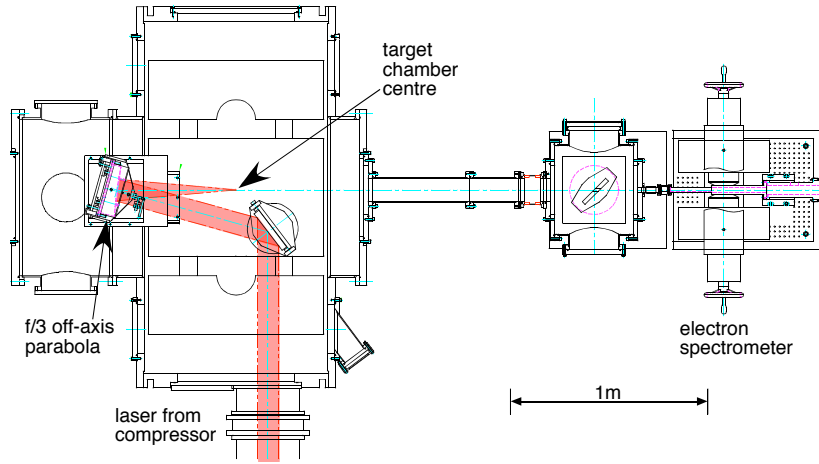


Figure 3.12: Sketch of the target chamber in TAW.

contrast ratio for VULCAN Petawatt is expected to be around 4×10^{-8} on a nanosecond timescale [158].

For one of the experiments described, the beam diameter was reduced in order to provide f/5 focusing and achieve a larger f-number, without changing the focal length. To achieve this an apodiser was placed in the beam at the end of the rod amplifier chain, just before the beam was sent through the disk chain.

3.3.4 Target Area West (TAW)

TAW has a 100 TW, sub-picosecond beam along with a 10 ps probe and up to 6 other beams each of 200 J in a nanosecond. The two short beams are created using the CPA technique described in section 3.3.1. For the experiment in TAW that is described in this thesis only the 100 TW and the probe beam were used. The pulse length of the main pulse was reduced to 640 fs, using the OPCPA technique described in 3.3.2.

In the front end of the standard 100 TW CPA line of TAW the ultrashort seed pulses are generated by a Time Bandwidth commercial oscillator (GLX-100) using Nd:Glass as the active medium. Mode locking is achieved by a semiconductor sat-

urable absorber mirror (SAM). The oscillator produces 1 nJ pulses of 170 fs duration at 80 MHz. The pulse is stretched using a standard refractive telescopic design operated in double pass mode, with a grating with 1740 lines per mm. The maximum bandwidth this system can pass at 1053 nm is 16 nm. Pre-amplification is performed by the 3-stage double-passed Nd:glass amplifiers, then propagated through one of the main VULCAN amplifier chains where it is initially amplified in a series of single-pass, flash-lamp pumped, rod-amplifiers to an output aperture of 45 mm. The final amplifier stages consist of a double-pass 108 mm aperture disc amplifier followed by a single-pass 150 mm aperture disc amplifier. The compressor in TAW consists of a single pass grating configuration with a 3.5 m separation, using gold-coated holographic gratings with 1740 lines per mm. A final pulse length of 1 ps can usually be achieved [159, 160].

In the experiment conducted in TAW which is described in this thesis, the pulse length was reduced to below 1 ps. This was achieved by increasing the amplified bandwidth, using the OPCPA front end of the PW laser as shown in figure 3.13. In this setup the oscillator was the Ti:Sapphire oscillator which is usually used for TAP. The pulse was then stretched by the TAW stretcher described above, and sent through the OPCPA amplifier, leading to an amplification to the mJ level. After that the pulse was then sent straight through the phosphate chain, skipping the silicate rod amplifiers [161].

The focusing geometry in TAW is a lot more flexible than in TAP, with the available parabolic mirrors ranging from $f/3$ to $f/20$. A schematic of the setup for the $f/3$ focusing geometry of the main pulse is shown in figure 3.12.

The second CPA beam can be re-compressed in air to 10 ps with 30 J. It is often used as a probe beam, and can also be converted to 2ω and 4ω . This has the advantage of being able to probe to higher densities, since the critical density for higher frequencies also increases.

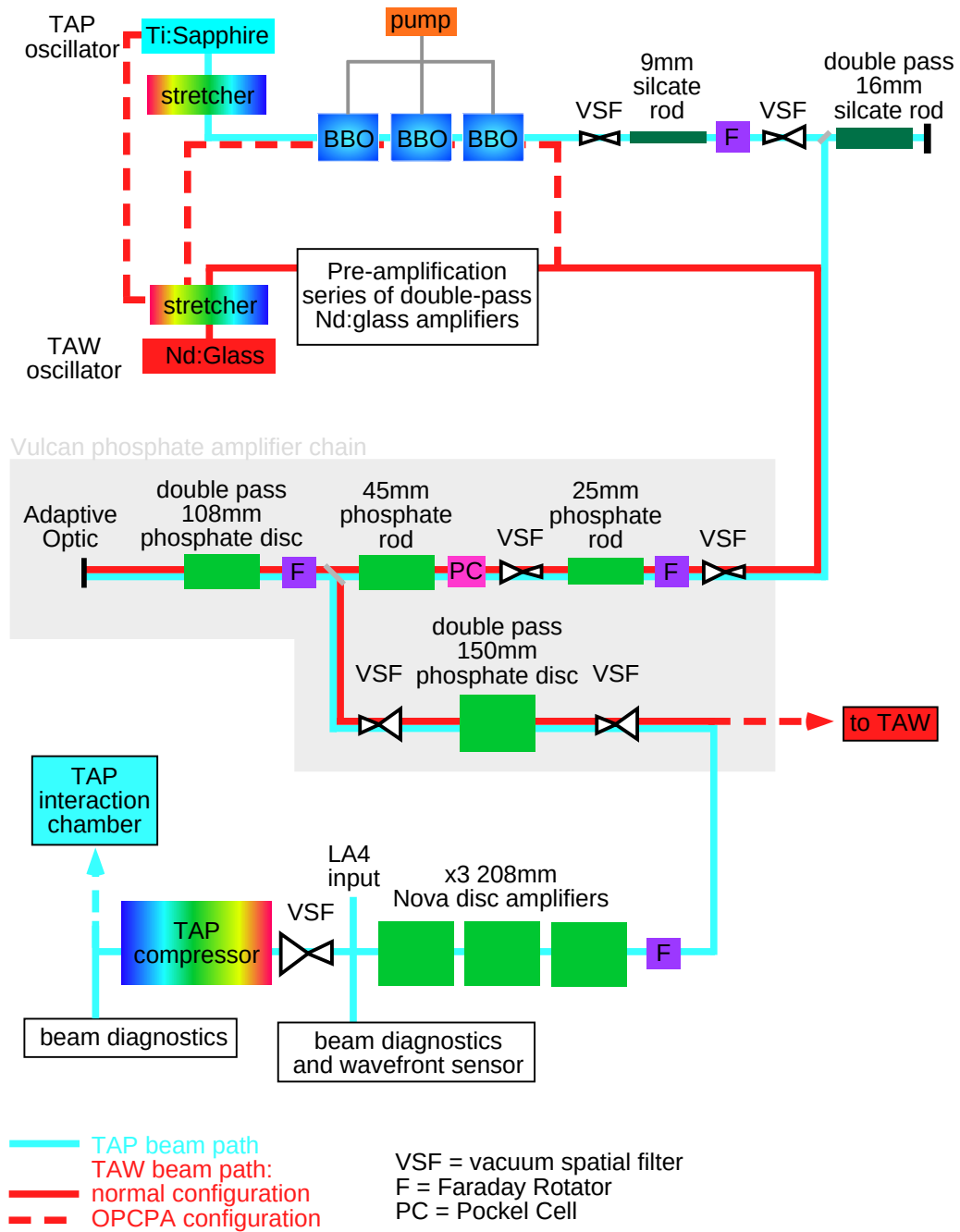


Figure 3.13: A schematic of the VULCAN Petawatt and 100 TW laser systems.

3.3.5 Parameters used in experiments

For the experiment in TAW in 2007, the OPCPA front end of the PW laser was used to reduce the pulse length from ~ 1 ps to (640 ± 120) fs. The focal spot size is determined by imaging the focal spot of a CW laser, co-linear to VULCAN, with a microscope objective. Confirmation that this is an accurate measurement of the focal spot size on a full power shot has been made by imaging 3ω emission from a solid target experiment [162], which gave a measurement of the focal spot that agreed to $\pm 0.5 \mu\text{m}$. The peak intensities given are the cycle averaged peak vacuum intensities for a gaussian pulse in space and time:

$$I_0 = \frac{8\mathcal{E}_{laser}}{\tau d^2} \left(\frac{\ln 2}{\pi} \right)^{\frac{3}{2}} = 0.829 \frac{\mathcal{E}_{laser}}{\tau d^2} \quad (3.9)$$

where $\mathcal{E}_{laser}$ is the laser energy, τ is the full width half maximum (FWHM) pulse length and d is the FWHM focal spot diameter (for a derivation, see appendix A.10).[‡] The intensities obtained with formula (3.9) are slightly higher than the values for the measured intensity in the focus, $I = 0.6 * \frac{4\mathcal{E}_{laser}}{\tau \pi d^2}$ [57], by about a factor of 9%. The laser parameters for the different experiments are summarised in table 3.3.

Underdense acceleration regime	run-II		run-I	$I < 10^{20} \text{ Wcm}^{-2}$
Overdense – angle of incidence		solid 40°	solid 45°	solid 8°
Pulse length (fs) f/3	570±175	517±103	630±120	642±124
Pulse length (fs) f/5	–	–	760±100	–
Pre-compression E_{max} (J)	694	726	595	119
Percentage onto target	55%	47%	56%	57%
Focal spot diameter (μm)	5	8	5	9.5
Maximum average power (TW)	670	660	529	106
Peak intensity (Wcm^{-2})	2.2×10^{21}	8.6×10^{20}	1.8×10^{21}	9.7×10^{19}
Peak a_0	43	27	38	8.9
Contrast ratio	10^{-7}	10^{-7}	10^{-7}	$\approx 10^{-6}$

Table 3.3: A summary of the laser parameters for each of the experiments.

[‡]Self-focusing in the interaction may enhance the intensity significantly.

Chapter 4

Experiments with underdense plasmas

This chapter presents experimental data on electron acceleration from underdense plasmas obtained using the VULCAN laser. With the long pulse length that VULCAN produces ($\gtrsim 500$ fs) electron densities as low as $\approx 10^{16} \text{ cm}^{-3}$ would be necessary, in order to satisfy the optimal condition for laser wakefield acceleration (LWFA) which requires that the laser pulse length equals half the plasma wavelength. At these densities the dephasing length would be on the order of tens of meters. With similar parameters, as used in the experiments discussed in this chapter, namely intensities $\geq 10^{20} \text{ Wcm}^{-2}$ ($a_0 \gg 1$), laser pulse lengths of approximately 700 fs, and densities between 10^{18} cm^{-3} to 10^{20} cm^{-3} , hot electrons have been observed and simulations suggest that direct laser acceleration (DLA) is the dominant acceleration mechanism [27, 119].

In this chapter this parameter range is further investigated. This investigation includes an intensity scan, as well as a change of the f-number in order to control filamentation seen at higher densities in [27].

The chapter is organised in the following way, the first section describes experiments in the Target–Area–Petawatt (TAP) using intensities between $1.2 \times 10^{20} \text{ Wcm}^{-2}$ and $2.2 \times 10^{21} \text{ Wcm}^{-2}$. It starts with a description of the experimental setups for two experimental runs. The electron spectra from these experiments are then dis-

cussed together, including intensity scans, density scans using different f-numbers and interaction length scans. The second section focuses on an experiment in Target–Area–West (TAW) of the VULCAN laser. Here different f-numbers with intensities ranging from $1.4 \times 10^{18} \text{ Wcm}^{-2}$ to $8.2 \times 10^{19} \text{ Wcm}^{-2}$ are used, and the electron spectra are compared. This is followed by 2D particle-in-cell (PIC) simulations looking at the effects of different laser intensities and f-numbers.

4.1 Experiments with intensities $> 10^{20} \text{ Wcm}^{-2}$

4.1.1 Experimental setups

Both experiments, described in this section, were performed using the VULCAN Petawatt Laser. The central wavelength of VULCAN is 1054 nm. For this wavelength the non-relativistic critical electron density is $n_c = 1 \times 10^{21} \text{ cm}^{-3}$. The targets were supersonic gas jets (see section 3.2), with different nozzle diameters (1, 2, 3 and 5 mm). By focusing the laser into the gas, underdense plasma is formed.

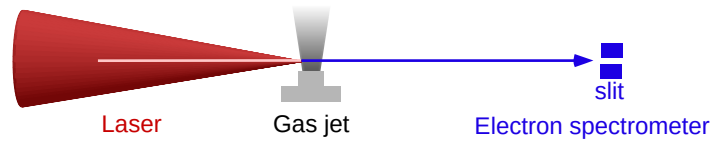


Figure 4.1: Sketch of experimental setup.

For the first run (run-I), the laser had a duration of (630 ± 120) fs and the laser energy on target ranged from 56 J to 336 J. A focal spot measurement, for the f/3 focusing geometry, at low power gave a FWHM of approximately $5 \mu\text{m}$. Leading to intensities ranging from $3.0 \times 10^{20} \text{ Wcm}^{-2}$ to $1.8 \times 10^{21} \text{ Wcm}^{-2}$. The accelerated electrons were measured in the direction of the laser propagation (i.e. at 0°). A 7.3 mm wide slit was placed in front of the entrance to the electron spectrometer to limit the angular acceptance and hence improve the energy resolution. The slit size was chosen as a compromise between the energy resolution of the electrons and signal level and size for the detection of x-rays emitted from the interaction, which was performed simultaneously (see [28]). In addition to the electron spectrometer,

electron beam divergence stacks were used to discern the electron beam divergence above a certain energy threshold. The stack was placed 10 cm behind the laser focus.

In the second run (run-II) a series of shots were taken at higher intensity, where the laser had a pulse duration of (570 ± 30) fs and energy of up to 485 J on target. Again a f/3 off-axis parabolic mirror focused the laser onto the edge of a gas jet. The FWHM of the focal spot was $5 \mu\text{m}$. This implies intensities up to $2.2 \times 10^{21} \text{ Wcm}^{-2}$. Here a 2 mm slit was placed in front of the entrance to the electron spectrometer, to limit the angular acceptance and hence improve the energy resolution.

The experimentally observed energy distributions are usually compared with a quasi-maxwellian function ($y = A \exp(-E^{MeV}/T_{\text{eff}}^{MeV})$) to define the effective temperature [26] or otherwise with a power-law function [120].

4.1.2 Intensity scans

In run-I, a laser energy scan was carried out at an electron density of $1.6 \times 10^{19} \text{ cm}^{-3}$, using a 2 mm nozzle and f/3-focusing. Figure 4.2 shows the obtained electron spectra. Two distinctly different shapes can be distinguished. At low laser energy (here $a_0 = 16$), the spectrum is clearly non-maxwellian. Such spectra are regularly observed in LWFA, when trapping of electrons in a small volume of phase space is achieved. At higher a_0 , the spectra become maxwellian. DLA could have lead to thermalisation of initially wakefield accelerated electrons.

Figure 4.3 shows the effective temperatures of the spectra in figure 4.2 as a function of intensity. The open symbols indicate the temperature obtained from a manual fit, which weighs all electron energies in the spectra equally. Whereas a simple maxwellian/exponential fit (solid symbols) puts more weight on the lower electron energies. This is due to the fact that more data points are available at these lower energies because of the non-linear dispersion of the magnet. The two curves in figure 4.3 are only fitted to the data points at higher a_0 , and show a linear fit ($\propto a_0$) and the scaling of the resonance energy $\gamma_e \propto (a_0 + a_0^3/2)$ (see appendix A.7). The temperatures for the higher laser energies seem to follow these scalings.

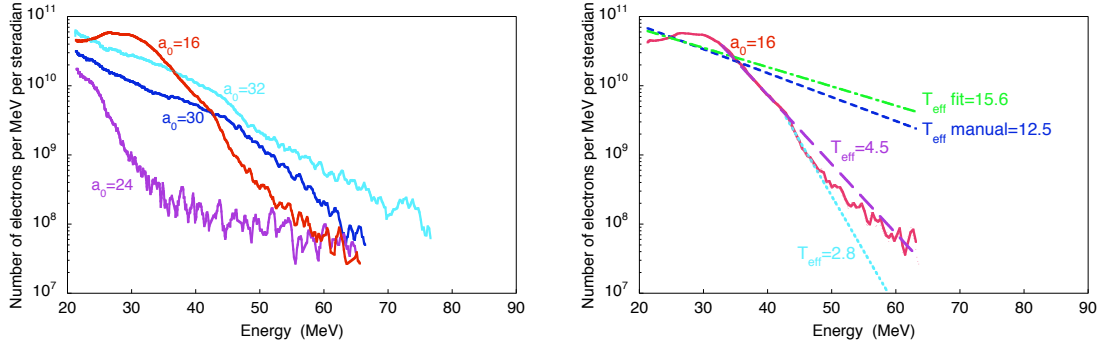


Figure 4.2: Intensity scan with f/3-focusing, using a 2 mm diameter nozzle with $n_e = 1.6 \times 10^{19} \text{ cm}^{-3}$.
 $a_0=16$ corresponds to $\mathcal{E}_L=56 \text{ J}$, and $I = 3 \times 10^{20} \text{ Wcm}^{-2}$;
 $a_0=24$ corresponds to $\mathcal{E}_L=128 \text{ J}$, and $I = 6.8 \times 10^{20} \text{ Wcm}^{-2}$;
 $a_0=30$ corresponds to $\mathcal{E}_L=207 \text{ J}$, and $I = 1.1 \times 10^{21} \text{ Wcm}^{-2}$;
 $a_0=32$ corresponds to $\mathcal{E}_L=236 \text{ J}$, and $I = 1.24 \times 10^{21} \text{ Wcm}^{-2}$;
The figure on the right shows fits to different parts of the spectrum for $a_0=16$.

Though with this limited data, it is impossible to differentiate the two. However, for the non-maxwellian spectrum (for $a_0 = 16$), the temperature does not appear to scale in the same way. This implies that the acceleration could be different for this lower intensity. Another indicator of the non-maxwellian shape of the spectrum is the large error bar obtained using the manual fit method. The fits to the different parts of the spectrum for $a_0 = 16$ are shown on the right hand side of figure 4.2.

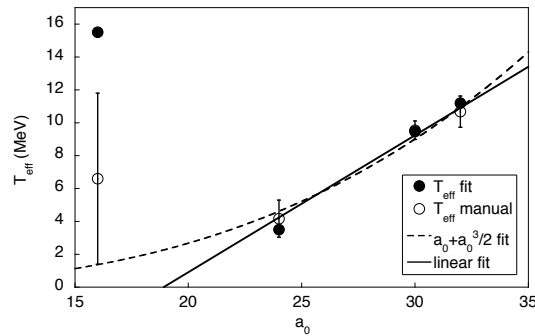


Figure 4.3: Scaling of the temperatures of the spectra in figure 4.2, including a linear fit (solid line) and a $a_0 + a_0^3/2$ fit (dashed line) for $a_0 > 20$.

This theory is supported by simultaneous x-ray measurements [28], which showed

that the x-ray radiation emitted in the forward direction decreased dramatically in number of photons with laser intensity. Taking into account the x-ray spectrum and number of photons [163] and comparing with the electron spectrum and number of electrons, it was inferred that the oscillation amplitude of the electrons increased with increasing laser intensity, indicating a transition from a wakefield dominated regime to a regime where electrons are accelerated by DLA. Here higher laser intensities drive larger transverse oscillations. For the case of the lowest a_0 , the acceleration is mostly axial, and the x-ray radiation becomes significantly less.

The transmitted laser spectra in figure 4.4 show satellites at the plasma frequency, which indicates SMLFA. The figure on the right hand side shows the scaling of the sideband intensities with a_0 . For the highest intensity, the weakest satellites are observed, and the lowest intensity $a_0 = 16$ shows the most energy in the satellites. This intensity dependence further supports the greater importance of SMLFA at lower a_0 .

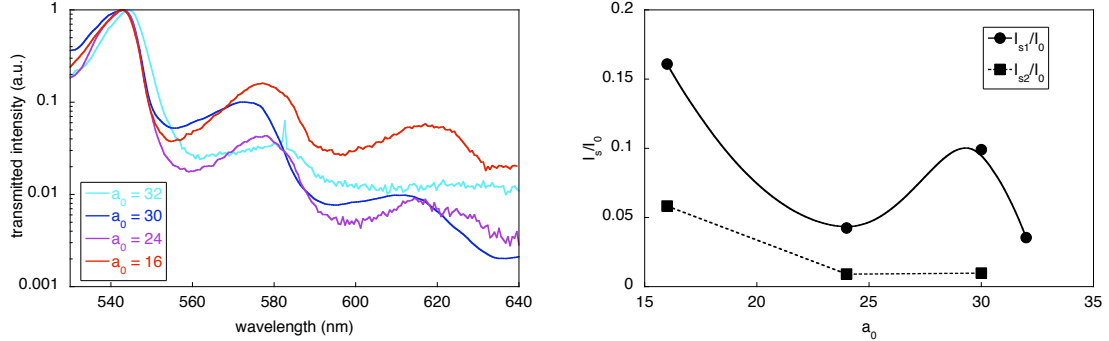


Figure 4.4: Spectra of the transmitted laser for different laser intensities, corresponding to the electron spectra in figure 4.2. And the scaling of the sideband intensities with a_0 .

Figure 4.5 shows another energy scan (taken on run-II), this time in hydrogen at an electron density of $2 \times 10^{19} \text{ cm}^{-3}$, again from a 2 mm nozzle, and using the f/3 focusing geometry. Note that here the spectra are plotted on a *log-log* scale, in order to show the power law shape of the spectra. The number of electrons per steradian increases with increasing laser energy. For the highest laser energy the signal descends into the background at 270 MeV. It is notable that the electron

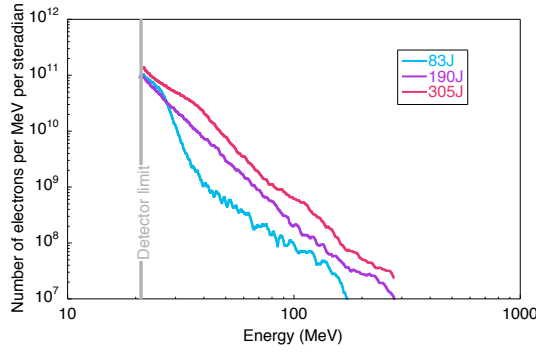


Figure 4.5: Energy scan, with hydrogen and a 2 mm nozzle at $2 \times 10^{19} \text{ cm}^{-3}$. The electron spectra were taken at the corresponding energies on target shown in the legend.

83 J corresponds to $a_0=18$, and $I = 4.1 \times 10^{20} \text{ Wcm}^{-2}$;

190 J corresponds to $a_0=28$, and $I = 9.4 \times 10^{20} \text{ Wcm}^{-2}$;

305 J corresponds to $a_0=35$, and $I = 1.5 \times 10^{21} \text{ Wcm}^{-2}$;

energies measured in this experiment are greater than the linear dephasing energy ($2 \frac{n_c}{n_e} m_e c^2 \approx 50 \text{ MeV}$), but smaller than the non-linear one, see section 2.2.3.6. For the 83 J shot, the spectrum looks more like a two temperature spectrum. This two-temperature shape of the spectrum may indicate a two-stage acceleration.

The shape of the spectra, obtained at higher intensities, is better fitted by a power law, rather than an exponential function, which can be seen by the linear behaviour on the *log-log* plot. Therefore it is harder to estimate an effective temperature for these spectra. The temperatures shown in figure 4.6 (open squares) were obtained by using the manual as well as the exponential fits, and taking the mean value. The fits in figure 4.6 are linear fits ($T_{\text{eff}} \propto a_0$) to the data from run-I for $a_0 > 20$, solid line, and for the data from both runs combined, again for $a_0 > 20$, dash-dotted line, as well as a $a_0 + a_0^3/2$ fit, dashed line, for $a_0 > 20$.

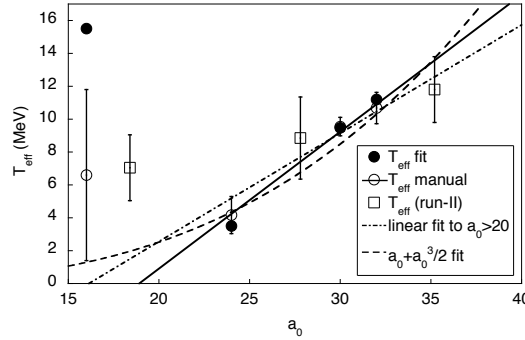


Figure 4.6: Same as figure 4.3 with the added estimated temperatures for the spectra in figure 4.5 (open squares). The lines are fits to the data for $a_0 > 20$, the solid line is a linear fit to the data from run-I, the dash-dotted line is a linear fit to the data from both runs combined and the dashed line is a $a_0 + a_0^3/2$ fit.

4.1.3 f-number comparison and density scans

This section investigates the effect of electron density on the electron acceleration, and compares different f-numbers.

The laser was focused with either f/3 or f/5 focusing geometry. For the f/3 focusing geometry, the laser energy on target ranged from 85 J to 336 J, and for the f/5-focusing, it was 85 J. For the two different f-numbers, f/3 and f/5, the laser had a duration of (630 ± 120) fs and (760 ± 100) fs respectively. This change in pulse length was due to a loss of bandwidth for the apodised (f/5) beam. Since the beam is spatially chirped, it loses bandwidth, due to the apodising, which results in a longer compressed pulse. A focal spot measurement for the case of f/3, at low power gave a FWHM of approximately $5 \mu\text{m}$. A diffraction limited spot* would have a FWHM of $3.1 \mu\text{m}$ and $5.4 \mu\text{m}$ for f/3 and f/5 focusing respectively, which makes the experimentally measured spot for f/3 1.6 times diffraction limited. From this we can infer a FWHM of about $9 \mu\text{m}$ for the case of f/5. The intensities on target for f/3 and f/5 can be estimated to be $4.5 \times 10^{20} \text{ Wcm}^{-2}$ and $1.2 \times 10^{20} \text{ Wcm}^{-2}$ respectively for 85 J on target. The spot size also influences the Rayleigh length ($z_R \approx 2f_{\text{number}}d_{\text{FWHM}}$, see appendix A.6), which then becomes approximately $30 \mu\text{m}$

*The first zero of an Airy function appears at $1.22f\lambda$ from axis. This distance is the same as 1.18 FWHM.

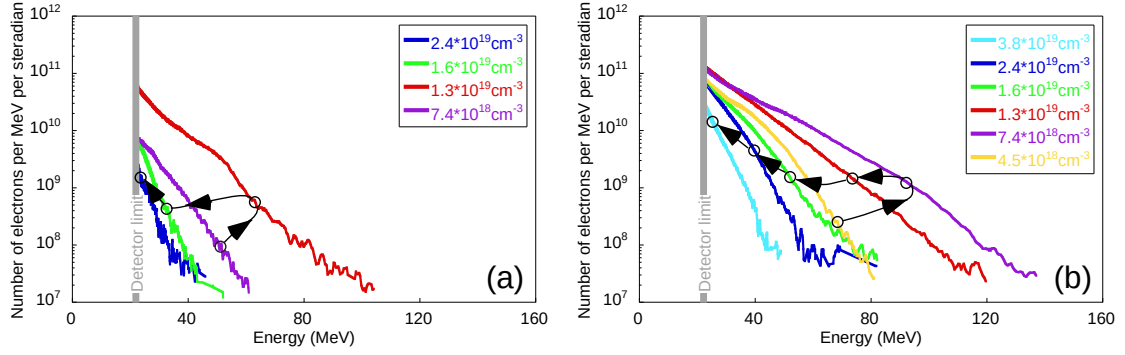


Figure 4.7: Electron spectra from 5 mm nozzle, for density scans using different f-numbers, namely (a) f/3-focusing, and (b) f/5-focusing. $\mathcal{E}_L = 85$ J. The arrows in both figures follow the increasing densities.

for the case of f/3 and $90 \mu\text{m}$ for the case of f/5.

Figure 4.7 shows electron spectra for density scans taken with f/3 (a) and f/5 (b) focusing. In the electron spectra the number of electrons per MeV per steradian (log scale) is plotted versus the energy of the electrons. The different colours correspond to different electron densities, and the effective temperatures of the electrons can be obtained by exponential fits to the spectra.

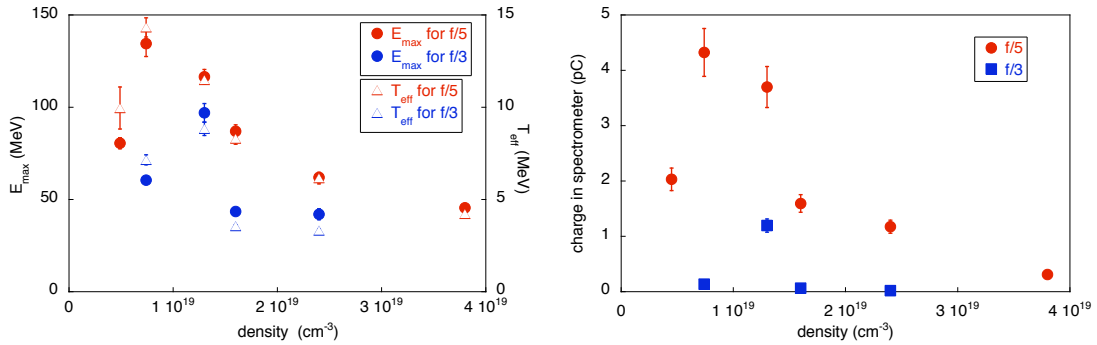


Figure 4.8: f/3 (blue) and f/5 (red) comparison for maximum electron energies (dots) and effective temperatures (triangles). The plot on the right hand side shows the measured charge as a function of density. The laser energy on target was $\mathcal{E}_L = 85$ J, using a 5 mm nozzle.

The obtained values for the effective temperature as well as the observed maximum energies are plotted against electron density in figure 4.8. Surprisingly the effective temperatures, maximum energies and charge are higher for the f/5-focusing,

even though it features a lower laser intensity. The maximum energy is about a factor of ten larger than the effective temperature, but both follow the same trend with density. One can easily see that both density scans show an optimum density (large T_{eff} , E_{max} and measured charge). A similar variation of energy with density has been observed, but not discussed before in [27, 142].

First the total charge and effective temperature increase with decreasing density. At densities below the optimum, the initial trapping and LWFA of the electrons is unlikely and hence less electrons will be injected into the channel. The acceleration to different maximum energies could be due to shorter acceleration lengths for higher densities [164]. This density scaling can be found for both f-numbers, but the optimum density decreases for larger f-number.

The blow-out regime, or cavitation, occurs when $I_{18}\lambda_\mu^2 > n_{18}d_\mu^2/20$ [78], where n_{18} is the electron density in 10^{18} cm^{-3} and d_μ is the FWHM of the focal spot in μm . This condition is met for all densities in the experiment, apart for the shot at $3.8 \times 10^{19} \text{ cm}^{-3}$ using f/5-focusing and may be even met then, due to self-focusing. For the blow-out regime: $\omega_\beta = \frac{\omega_p}{\sqrt{2\gamma_e}}$ [28]. When $\omega_\beta \simeq \frac{1}{2\langle\gamma\rangle} \frac{n_e}{n_c}$ (see equation A.75) resonance between the betatron oscillations and the laser field occurs (assuming $v_{ph} \approx c$). Therefore the energy at which betatron resonance occurs scales as $\gamma_e \propto 1/n_e$. The scaling of the temperature with density for the case of f/5-focusing is plotted in figure 4.9. The non-linear dephasing energy for the wakefield accelerator also scales as $1/n_e$, but would be higher than the observed maximum energies. A reduction of a_0 by a factor of 5, for the case of f/5-focusing, would be necessary to match the measured energies, but the lowest densities would not scale, either. Furthermore, one would expect higher energies for the higher a_0 of the f/3-focusing.

If one only takes the intensity dependence for DLA into account (see section 4.1.2, figure 4.2), this would suggest that f/3-focusing, giving a higher a_0 , should lead to better electron acceleration than the case of f/5-focusing. However, in the case of f/3 we observe lower temperatures, maximum energies and charge, than in the case of f/5. In other words, the acceleration of the electrons is more efficient in the case of f/5 over the range of densities studied here.

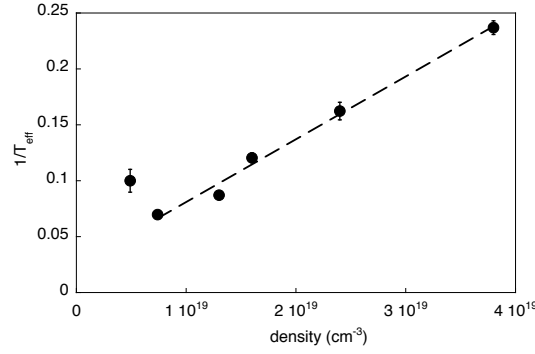


Figure 4.9: $1/T_{\text{eff}}$ for f/5-focusing plotted against density. The dashed line is a linear fit to the data points excluding the one at $4.4 \times 10^{18} \text{ cm}^{-3}$.

A possible explanation could be that the f/5 focusing geometry features a larger focal spot and longer pulse length, but lower a_0 . γ_e , the energy an electron has to reach before it can gain energy from the betatron resonance, is also dependent on a_0 , $\gamma_e \propto (a_0 + a_0^3/2)$ (see appendix A.7). Therefore, γ_e is almost a factor of 10 larger for f/3 than it is for f/5, making it harder for electrons to meet the resonance condition in the case of f/3. This means that the larger f-number can more easily accelerate electrons in the DLA. The enlarged focal spot will lead to an increase in channel size, which could support more oscillating electrons in case of f/5 explaining the increase of the accelerated charge from f/3 to f/5.

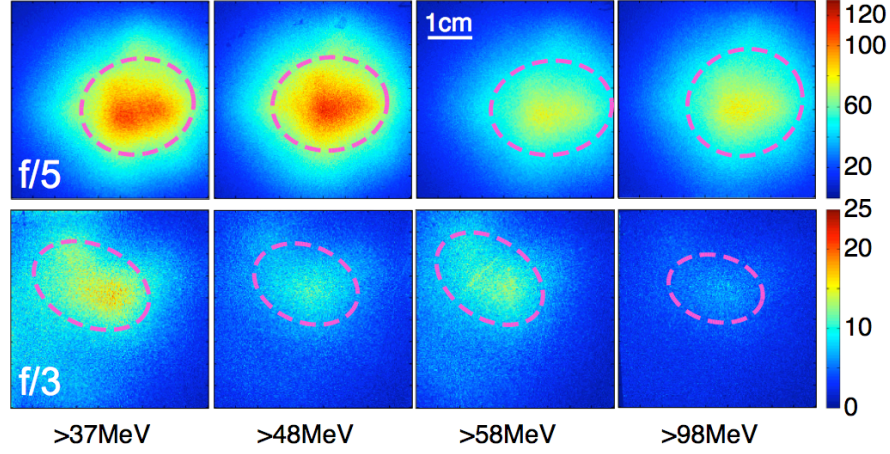


Figure 4.10: Stack data for shots taken at a density of $1.3 \times 10^{19} \text{ cm}^{-3}$, for f/5-focusing (top) and f/3-focusing (bottom). The energy threshold increases from left to right. The colour bar indicates the number of electrons per pixel. The contours (dashed pink) show the FWHM of the electron beam divergence.

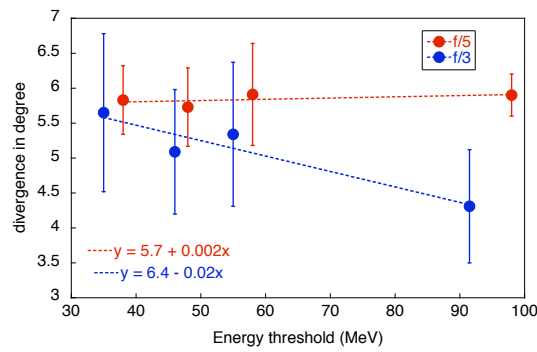


Figure 4.11: Beam half-angle divergence at the FWHM from the stack data, in figure 4.10, at the different energy thresholds, f/5 (red) and f/3 (blue). The dotted lines are linear fits to the points.

Figure 4.10 shows the electron stack data as a number of electrons per pixel. The top row is the imaging plate detector data for the shot using f/5 focusing geometry, and the bottom row for the f/3-focusing. The shots were taken at an electron density of $1.3 \times 10^{19} \text{ cm}^{-3}$, using a 5 mm nozzle. As we go through the stack, the signal in the case of the f/3-focusing drops significantly with increasing energy (low temperatures) whereas in the case of the f/5-focusing the signal drops slower (higher temperature). The stack data and the data from the spectrometer show the same trends, namely higher number of electrons are detected in the case of f/5, and the temperature in the case of f/5 is higher as well. Furthermore, the energy up to which we observe electrons is higher for the case of f/5.

Figure 4.11 shows the divergence angles for the FWHM (pink contours in figure 4.10) of the detected electron beam from the stack data. This is plotted against the energy threshold of the detectors position. It can be observed that in the front of the stack, the beam divergence is similar for f/3 and f/5. As we go further back in the stack, the divergence for f/5 remains approximately constant, whereas the divergence for f/3 decreases gently. This suggests that in the case of f/5-focusing the electrons pick up radial as well as longitudinal momentum, supporting the assumption that higher energies are accelerated by DLA.

Using the divergence angles of the electron beams for the given density of $1.3 \times 10^{19} \text{ cm}^{-3}$, in conjunction with the obtained spectra gives the total charge in the electron beam as 1.3 nC for the case of f/3 and 5.7 nC for f/5-focusing. This gives an energy conversion efficiency of 0.02% for the f/3-focusing and 0.08% for the f/5-focusing. This agrees with the amount of electrons detected using the stacks, indicating that the majority of the signal is due to electrons and not secondary particles.

A divergence measurement was also conducted on run-II with the f/3 focusing geometry. Figure 4.12 shows the electron beam divergence at three different positions in a stack positioned 94 mm behind focus. The data for figure 4.12 were taken at the same parameters as those for figure 4.13. The pink contours indicate the FWHM of the beam divergence.

The shape of the electron spectra observed in run-II was not purely maxwellian,

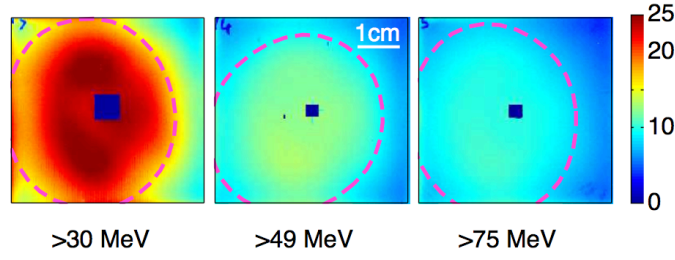


Figure 4.12: Imaging plate data of a stack for a shot taken with the same parameters as for figure 4.13. It shows the electron beam divergence (FWHM pink contours) at energy thresholds corresponding to ≥ 30 MeV, ≥ 49 MeV and ≥ 75 MeV.

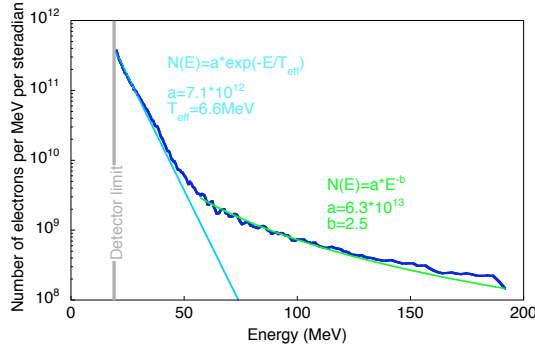


Figure 4.13: Electron spectrum taken using a 2 mm nozzle at a density of $4 \times 10^{19} \text{ cm}^{-3}$ with (light blue) exponential fit and (green) power law fit to energies ≥ 75 MeV.

but had an increased number of electrons in a higher energy tail. Figure 4.13 shows a typical electron spectrum for this experiment taken at an electron density of $4 \times 10^{19} \text{ cm}^{-3}$. The number of electrons per MeV per steradian (log scale) is plotted versus the energy of the electrons in MeV (linear scale). The blue fit is an exponential fit, and the lower energy part of the spectrum shows maxwellian behaviour. From the exponential fit it is possible to determine the effective temperature as 11.5 MeV. The green curve is a power law fit to the spectrum, starting at an electron energy of 75 MeV. The parameters for the two different fits are displayed in the graph.

Note that the maximum electron energy detected in the stack is 75 MeV, which is the energy after which the hotter tail in the spectrum in figure 4.13 starts. Since for

E threshold (MeV)	$\theta(^{\circ})$
30	11.1 ± 0.7
49	10.3 ± 0.6
75	10.0 ± 0.8

Table 4.1: The divergence half angle of the electron beam from the FWHM width, as shown in figure 4.12.

the stack shot, the signal is taking up most of the area of the detector, it is virtually impossible to discern a real background. Therefore, the widths of the FWHM might be smaller in reality, but the differences between the different divergences at the different positions in the stack should be the same for the different energies. The values of the beam divergence half angles, for the data shown in figure 4.12, are summarised in table 4.1. It can be seen, that the beam divergence does not change much with the energy of the electrons. Using the obtained divergence in conjunction with the spectrum in figure 4.13 yields a total charge of 10.3 nC. The charge in this run is higher, compared to run-I, because of the higher laser intensity (in run-II), and the energy conversion efficiency is comparable to the one for the f/3-focusing in run-I, with 0.023%.

4.1.4 Interaction length scans

In this section, the effect of interaction length is investigated for two f-numbers and intensities. An interaction length scan from run-I, using f/5-focusing is followed by one from run-II using f/3-focusing.

In figure 4.14 the electron spectra variation as a function of interaction length is shown. The scan was conducted at the same density, $1.6 \times 10^{19} \text{ cm}^{-3}$, as the intensity scan in figure 4.2. A laser energy of 90 J was focused using the f/5 focusing geometry, leading to an intensity of $1.1 \times 10^{20} \text{ Wcm}^{-2}$ ($a_0 = 9.5$). For the short nozzles (1 mm and 2 mm) we observe strongly non-maxwellian spectra. Two mechanisms can lead to the thermalisation of the electron spectra with interaction length. Assuming LWFA to be the dominant acceleration mechanism, acceleration in excess of

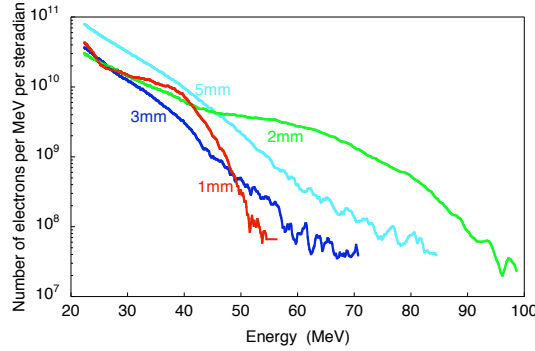


Figure 4.14: Interaction length scan using $f/5$ focusing geometry with nozzles of diameters, 5 mm, 3 mm, 2 mm and 1 mm. The density is $1.6 \times 10^{19} \text{ cm}^{-3}$ and the energy on target is 90 J.

the dephasing length will lead to an increase of the electron energy spread, due to the decelerating part of the wakefield potential. Therefore, with increasing interaction length, the thermalisation of the pre-accelerated electrons increases. The second mechanism that can lead to thermalisation of the electron spectra is DLA. Considering pre-accelerated electrons that do not follow a quasi-maxwellian distribution, as long as the volume they occupy in phase space is not too small, such electrons will have a random phase with respect to the laser. In [27] it was shown that even if the electrons start within a fraction of one λ , they still follow very different trajectories. Only a few of these electrons will fulfill the betatron resonance condition. For some of the electrons, stochastic heating, (section 2.2.3.7), can further increase the electron energy spread.

Figure 4.15 shows the electron spectra for three different nozzle diameters for run-II. Here the laser was focused using a $f/3$ focusing geometry, to obtain higher intensities. The laser energy and electron density were kept constant, at (330 ± 14) J and $2 \times 10^{19} \text{ cm}^{-3}$ respectively, as the nozzle diameter was changed, and the intensity was $\approx 1.6 \times 10^{21} \text{ Wcm}^{-2}$. However, both the intensity and the laser energy were higher than for the interaction length scan in run-I, discussed above. Due to the high signal level of the x-ray peak on the imaging plates, the maximal energy that can be resolved here is around 240 MeV. The spectra for the different interaction lengths look very similar and the temperatures of the lower energy part fluctuate

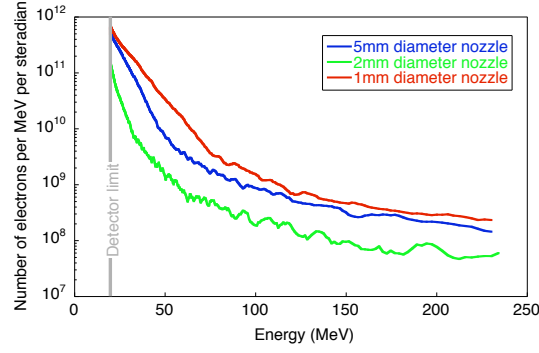


Figure 4.15: Comparison of the different nozzle sizes (interaction lengths), with otherwise similar parameters, at an electron density of $2 \times 10^{19} \text{ cm}^{-3}$, using f/3 focusing geometry.

around $7 \pm 2 \text{ MeV}$, whereas the ones for the higher energy tails are around 30 MeV. The overall shape of the spectra does not change very much, which indicates that most of the acceleration happens within the first mm.

4.2 Experiments with intensities $< 10^{20} \text{ Wcm}^{-2}$

Since the change from f/3– to f/5–focusing already showed a major improvement, going to even larger f–numbers would be interesting. However, this is not possible in TAP, and therefore the experiment described in this section was carried out in the Target–Area–West of the VULCAN laser.

The focusing geometry was changed from f/3 to f/5, f/10 and f/20, where the energy of the laser was the same for f/3 and f/20, but decreased for f/5 and f/10. The effects of these different conditions are discussed.

4.2.1 Experimental setup

The laser was focused with either f/3, f/5, f/10 or f/20 focusing geometry. The laser had a duration of (642 ± 124) fs. A focal spot measurement, for the case of f/3, at low power gave a FWHM of $(9.5 \pm 0.9) \mu\text{m}$. From this we can infer a FWHM of approximately $16 \mu\text{m}$ and $32 \mu\text{m}$, for the cases of f/5 and f/10 respectively. To achieve these geometries the f/3 focusing parabola is used and the beam diameter is reduced (apodised) to achieve the larger f–numbers. This leads to a reduction in the laser energy. The f/20–focusing on the contrary was achieved using a different off–axis parabola. The various laser parameters for the different focusing geometries of this experiment are summarised in table 4.2.

Parameters	f/3	f/5	f/10	f/20
Focal spot diameter (μm)	9.5 ± 0.9	16	32	38.8 ± 1.8
Energy on target (J)	57	46	12	57
Peak intensity (Wcm^{-2})	8.2×10^{19}	2.3×10^{19}	1.4×10^{18}	4.8×10^{18}
Peak a_0	8.2	4.4	1.1	2

Table 4.2: A summary of the laser parameters for the investigation of different focusing geometries.

Supersonic gas jet nozzles, with diameters of 2, 3, 5 and 10 mm, created a gas column with a sharp interface between the gas and the surrounding vacuum.

The electrons were measured in the direction of the laser propagation at 0° . An 8 mm wide slit was placed in front of the entrance to the electron spectrometer, to limit the angular acceptance and hence improve the energy resolution.

4.2.2 f-number comparisons

Figure 4.16 shows the electron spectra for three different f-numbers (f/3, f/5 and f/10), using a 3 mm diameter nozzle at a density of $3 \times 10^{19} \text{ cm}^{-3}$. Even though the laser energy for f/10 is a lot lower than for f/3 or f/5, the electron temperature and the maximum electron energy, as well as the number of electrons is larger.

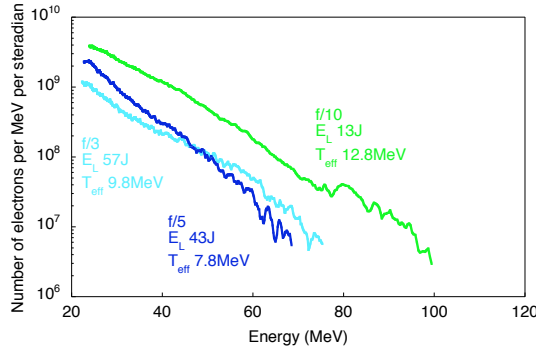


Figure 4.16: f-number comparison for a 3 mm nozzle at $3 \times 10^{19} \text{ cm}^{-3}$.

The expected dephasing energy (see equation (2.28)) for the f/10-focusing is around 41 MeV. This is about a factor of two smaller than the actually observed maximum electron energy of 100 MeV. However, self-focusing to a slightly higher intensity of $1.6 \times 10^{18} \text{ Wcm}^{-2}$, which equals $a_0 = 2.4$ can explain these energies. The dephasing length of 0.2 mm is a lot shorter than the gas jet length of 3 mm for all focusing geometries. It is also shorter than the Rayleigh length for all f-numbers, apart from the f/3-focusing, for which the Rayleigh length is shorter than 0.2 mm.

Shorter f-numbers are less efficient for the electron acceleration despite the higher laser energy, as they produce lower electron energies and temperatures. This could be due to better self guiding and less filamentation for the f/10 focusing geometry, leading to a higher effective intensity throughout the interaction with the plasma.

In figure 4.17 electron spectra for a 2 mm nozzle, comparing $f/20$ and $f/3$ at $2.4 \times 10^{19} \text{ cm}^{-3}$, are plotted. Comparing the shot at $2.4 \times 10^{19} \text{ cm}^{-3}$ for the $f/3$ focusing geometry with the one in figure 4.16, shows that the shorter propagation length leads to an increase in electron number and energy whereas the temperature is almost the same. This may indicate that for shorter interaction length (here shorter by 1 mm) less electrons are lost due to the process of dephasing.

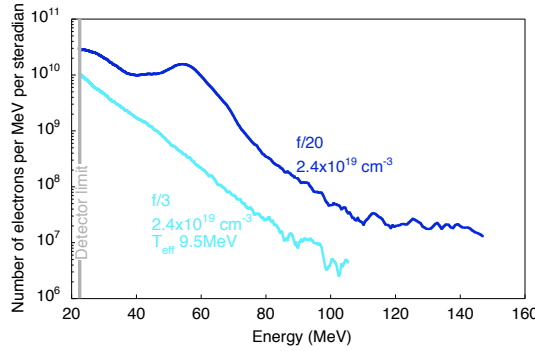


Figure 4.17: Electron spectra for 2 mm nozzle, comparing $f/20$ and $f/3$ at $2.4 \times 10^{19} \text{ cm}^{-3}$.

The spectrum for the $f/20$ focusing geometry is clearly non-maxwellian, and shows a bump, which indicates that dephasing has not thermalised the spectrum completely yet. This bunching is likely to be due to SMLWFA. For the longer f -numbers, as mentioned above, the self-guiding of the laser is expected to be more efficient.

Figure 4.18 shows the effective temperatures for three density scans using different f -numbers and a 3 mm diameter nozzle. Most of the spectra are well fitted by a maxwellian fit, indicating dephasing. However, the spectrum for $f/10$ -focusing at a density of $0.9 \times 10^{19} \text{ cm}^{-3}$ does not have an exponential shape. This can be seen in figure 4.19, which shows the density scan using $f/10$ -focusing.

In general, for all f -numbers, the temperature increases with increasing density and therefore does not show the previously observed n_e^{-1} dependency. However, again the largest f -number, produces the highest temperatures. For this experiment, most densities are not in the blow-out regime, with exception of the shots below $2 \times 10^{19} \text{ cm}^{-3}$ for the $f/3$ -focusing. Here $0.9 \times 10^{19} \text{ cm}^{-3}$ seems to be the optimum

density. This is roughly the same as in the higher intensity regime in the previous section.

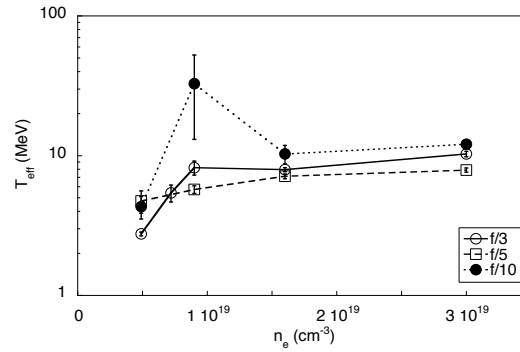


Figure 4.18: Effective temperatures for different densities and f-numbers.

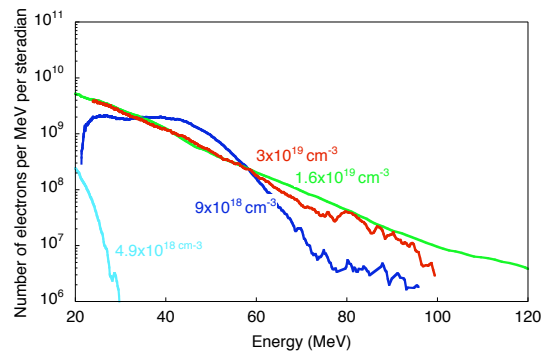


Figure 4.19: Electron spectra for a density scan using f/10-focusing, 3 mm nozzle, and 12 J on target.

4.2.3 Length scan

Figure 4.20 shows spectra for the f/10 focusing geometry at 2 different densities, comparing different interaction lengths. The spectra for the 10 mm nozzle for both densities show a decrease in the number of electrons, and maximum energy, as well as a more maxwellian shape of the spectra. This is in good agreement with the results from the interaction length scans in section 4.1.4, and can be attributed to a loss of energetic electrons due to dephasing.

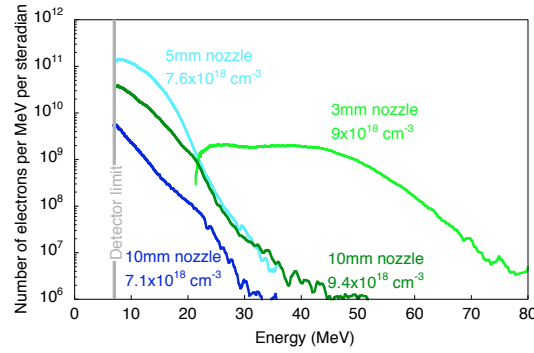


Figure 4.20: Electron spectra for f/10 using different nozzle diameters.

4.3 Particle-in-cell simulations

To gain better understanding of the experimental results numerical simulations were performed using OSIRIS. OSIRIS is a 3D-3V (3 dimensions in both space and velocity) collisionless particle-in-cell (PIC) code, written by the OSIRIS consortium of UCLA, IST (Portugal) and USC [165–167]. Running in 3 spatial dimensions is very computationally expensive and the resolution has to be decreased considerably. Therefore, for the simulations presented in this thesis, OSIRIS was run as a 2D-3V code at Imperial College London on the CX1 cluster. There are two ways to set up the simulation. One possibility is to use a box which co-propagates with the laser pulse at the speed of light in vacuum, c . The second way is to have a stationary box so that plasma dynamics can be observed. Simulations are initialised by preparing the input parameters of the density profile and the chosen laser parameters using the normalised units shown in table 4.3. The output files of the simulation are also in these normalised units.

unit	OSIRIS unit	conversion factors
time	$\frac{1}{\omega_0}$	0.559 fs
length	$\frac{c}{\omega_0}$	0.168 μm
velocity	c	
charge	e	
mass	m_e	
electron momentum	$m_e c$	0.511 MeV/c
density	$n_c = \frac{\epsilon_0 m_e \omega_0^2}{e^2}$	$1.0 \times 10^{21} \text{ cm}^{-3}$
electric field	$\frac{c m_e \omega_0}{e}$	3.05 TV/m
magnetic field	$\frac{m_e \omega_0}{e}$	10.2 kT

Table 4.3: The normalised units used in OSIRIS and useful conversion factors, for interactions at $\lambda_0 = 1.054 \mu\text{m}$.

4.3.1 Intensity scan

The simulations were conducted as a comparison to the intensity scan, described in section 4.1.2, where a_0 ranged between 16 to 32 at a density of $1.6 \times 10^{19} \text{ cm}^{-3}$. The simulation box is co-propagating with the laser pulse at the speed of light, c . A 500 fs long pulse with a gaussian temporal shape, and linear polarisation in the plane perpendicular to the simulation box, was focused to the edge of a 2 mm long, fully ionised helium plasma with a density of $1.6 \times 10^{19} \text{ cm}^{-3}$. The grid resolution of the $391 \mu\text{m}$ by $134 \mu\text{m}$ simulation box was 28 nm ($\approx \lambda_0/36$) in the propagation direction, z , and 84 nm ($\approx \lambda_0/12$) in the transverse direction (x -direction). The time step, the data was saved in, was 50 fs, and the transverse boundary conditions were Lindman open space [165, 168]. In order to perform intensity scans a_0 was changed and had the values of 5, 16, 24 and 30. The simulations were run for a focal spot size of FWHM $9.5 \mu\text{m}$.

Figure 4.21 summarises the electron density profile, electron momentum phase space, longitudinal and transverse electric fields, laser envelope as well as the perpendicular (azimuthal) magnetic field for three different intensities, after 1.41 mm of propagation. In general three parts in the simulation box can be distinguished. At the front of the laser pulse a “bubble” is formed. Behind the bubble, a transversely modulated part can be observed, which is followed by a channel.

For higher intensities the “bubble” at the front of the pulse is more pronounced, and the longitudinal electric field is a lot stronger. For the lowest intensity the injection in the first bubble is starting, but further back electrons have already been injected. At higher intensities, injection in the bubble has already happened earlier and the electrons are in the dephasing stage. The longitudinal electric field behind the bubble is unstructured and very weak, therefore the laser envelope was overlaid, using the yellow to dark red colour scale.

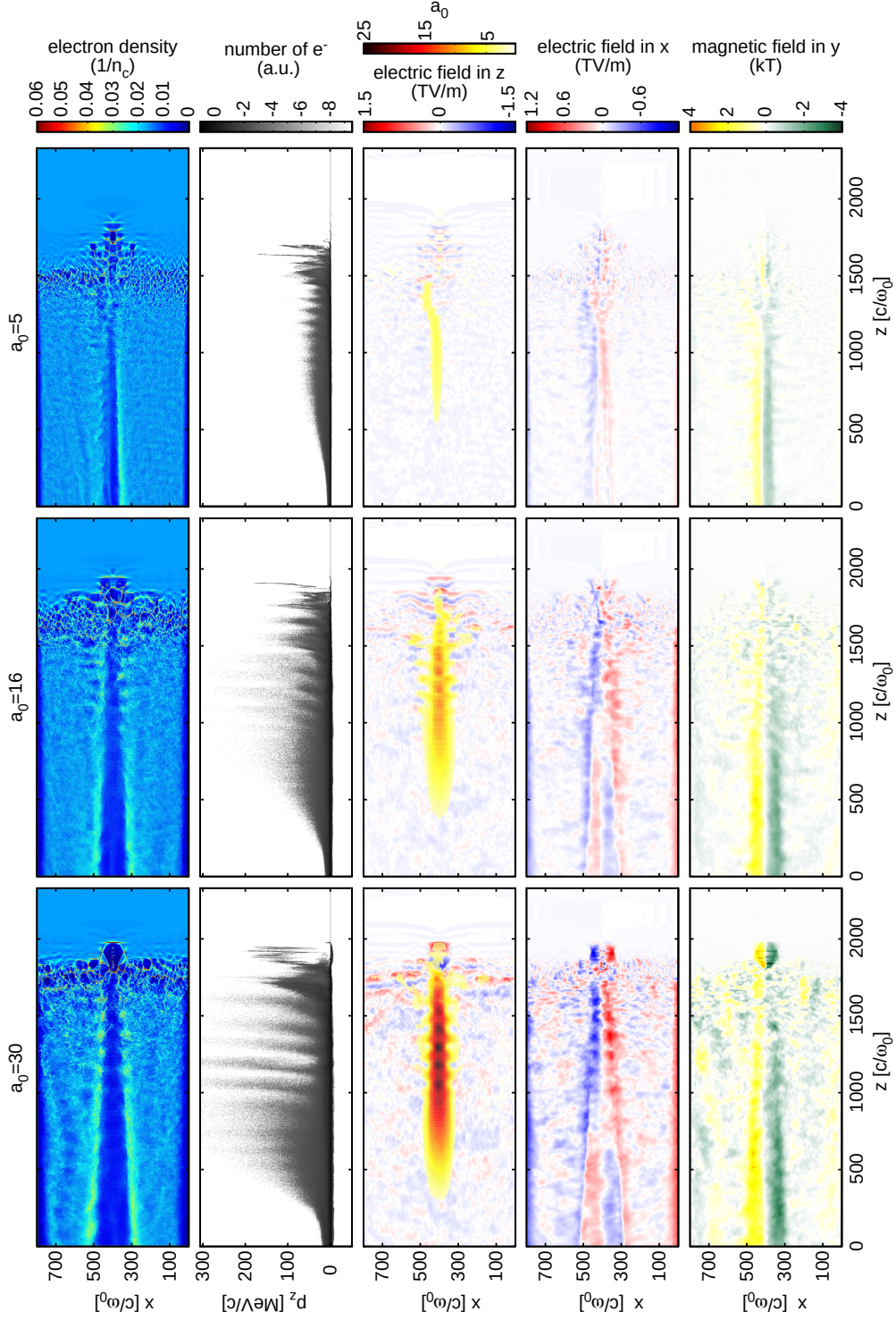


Figure 4.21: Plots of the electron density profile, electron phase space, longitudinal electric field and laser envelope, transverse electric field and perpendicular magnetic field for three different laser intensities, after 1.41 mm of propagation in helium with $n_e = 0.016n_c$.

$a_{0 \text{ max}}$	λ_{bunch}	$\lambda_{p \text{ rel}} = \sqrt{\langle \gamma \rangle} \lambda_p$	a'_0	$\lambda_{p \text{ rel}}(a'_0)$
5	$(11.7 \pm 0.9) \mu\text{m}$	$16 \mu\text{m}$	1.7	$10.5 \mu\text{m}$
16	$(17 \pm 2) \mu\text{m}$	$28 \mu\text{m}$	6.5	$18.1 \mu\text{m}$
24	$(20 \pm 2) \mu\text{m}$	$31.5 \mu\text{m}$	9.2	$21.4 \mu\text{m}$
30	$(23 \pm 2) \mu\text{m}$	$38.5 \mu\text{m}$	12.2	$24.6 \mu\text{m}$

Table 4.4: Relativistic plasma wavelengths, and bunching wavelengths from simulations.

The bunching of the electrons as seen in the phase space diagrams occurs at different wavelengths for different intensities. Table 4.4 lists these bunching wavelengths from the simulations (λ_{bunch}) and the relativistically corrected plasma wavelengths ($\lambda_{p \text{ rel}} = \sqrt{\langle \gamma \rangle} \lambda_p$), as well as the “average” measured a_0 in the relevant areas (a'_0) and the corresponding relativistic plasma wavelength $\lambda_{p \text{ rel}}(a'_0)$. This wavelength is close to the observed one. Figure 4.22 shows a plot of the values listed in table 4.4. The bunching at roughly the relativistic plasma frequency, which can also be seen in the laser envelope for higher a_0 ($= 16, 24$ and 30), and in the transmitted laser spectra, see figure 4.23, indicates SMLWFA. The intensity of the satellites decreases for increasing a_0 , similar to the experimental optical spectra in figure 4.4.

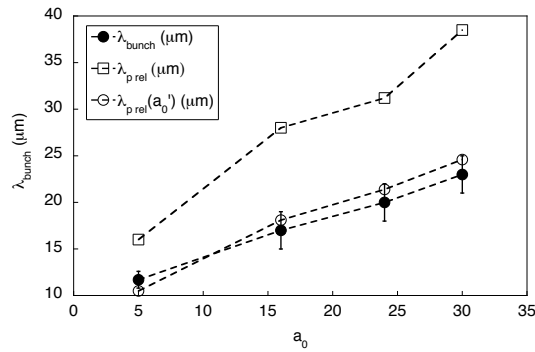


Figure 4.22: Bunch wavelength of electrons for different a_0 . λ_{bunch} is the observed bunching wavelength in the simulation, $\lambda_{p \text{ rel}}$ is the expected, relativistically corrected plasma wavelength and $\lambda_{p \text{ rel}}(a'_0)$ is the relativistically corrected plasma wavelength using the average a_0 from the simulation.

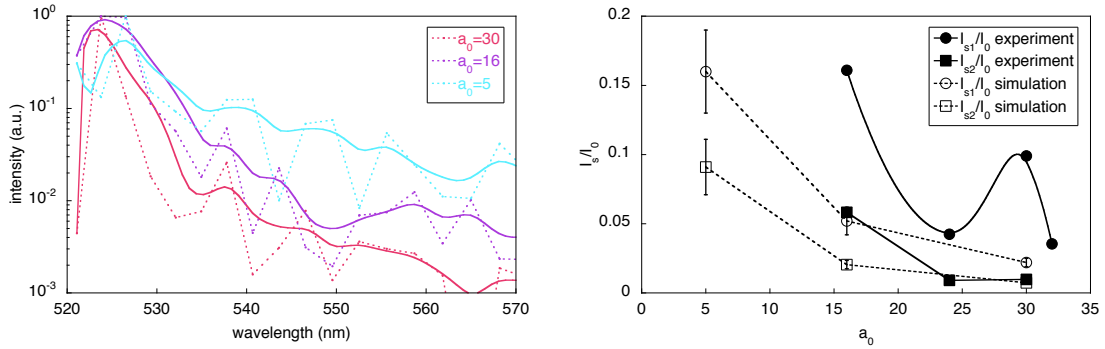


Figure 4.23: Optical spectra from the simulations, after 1.41 mm of propagation for different a_0 . The dashed lines follow the data, and the solid lines are a smoother representation. The figure on the right hand side shows the scaling of the sideband intensities with a_0 for both simulation (open symbols) and experiment (solid symbols).

In the third part of the simulation box, channel formation is observed. Here the modulation disappears, and the clear radial electric field and the azimuthal magnetic field, for the higher a_0 , both serve to trap hot electrons in the channel. A typical magnetic field strength of 2 kT (or 20 MG) and electric field strength of about 0.4 TV/m, which scale with a_0 , can be observed, see figure 4.24. Therefore $\nabla \mathbf{F} \propto a_0$ leads to a scaling of $\gamma_e \propto a_0 + a_0^3/2$ (see appendix A.7). This decrease in channel strength for lower intensity can also be seen in the phase space diagrams (2nd row in figure 4.21), where the energy the electrons are accelerated to in the trailing channel is lower than for high a_0 . The channel diameter also changes for different intensities and is largest for high a_0 , therefore supporting higher oscillation amplitudes.

Figure 4.25 shows the FWHM of the channels in the simulations for different a_0 , as well as the expected diameter of the channel in the blow out regime, for a simplified, triangular transverse intensity gradient, $2r_B = 4\sqrt{a_0}c/\omega_p$. This underestimates the channel diameter slightly, by approximately a factor of 1.5.

The temperatures of the phase spaces in figure 4.21 are plotted in figure 4.26. They do follow the linear scaling $T_{\text{eff}} \propto a_0$ like the experimental data, even for the lowest a_0 .

In figure 4.21 three different electron acceleration regimes are observed at the

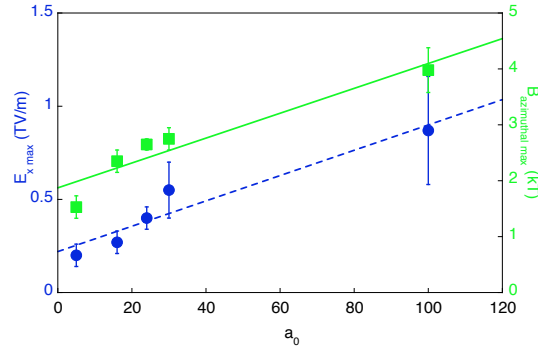


Figure 4.24: Dependence of the maximum transverse electric field and azimuthal magnetic field on a_0 .

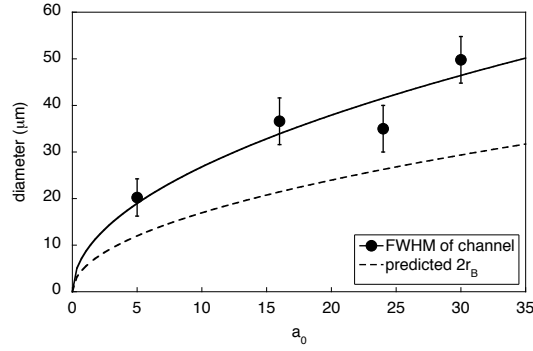


Figure 4.25: FWHM of the channels in the simulations for different a_0 . The expected $r_B = 2\sqrt{a_0}c/\omega_p$ is for a simplified model of a triangular transverse intensity shape.

same time: Bubble regime, SMLWFA and DLA. This is in partial agreement with the experimental data. The achieved electron energies increase with increasing a_0 , and for the lowest a_0 the SMLWFA seems to be the dominant mechanism, as DLA is less important.

Looking at the phase space for the three different intensities, after the laser has left the 2 mm long plasma, further enhances the last point. Figures 4.27 and 4.28 show clearly that for the lowest a_0 the maximum energy gained by the electrons is approximately the dephasing energy ($2m_e c^2 n_c / n_e = 64$ MeV). Whereas for higher a_0 , DLA plays a major role in accelerating the electrons further.

The electron spectra from the simulations however yield significantly higher tem-

peratures, and the electrons are accelerated to higher energies than in the experiment. This could be due to the fact that in the experiment, the pre-pulse of the laser ionises the gas, and can also pre-form a channel, so that the density, that the main part of the laser interacts with, is lower than the value used in the simulation. Other possible reasons for the difference between experiment and simulations could be numerical dispersion, or just the ideal nature of the simulated pulses.

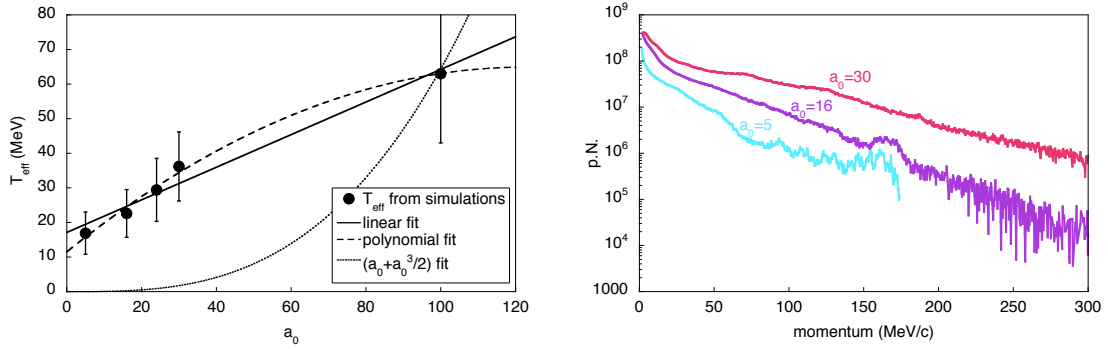


Figure 4.26: Temperatures after 1.41 mm propagation for different intensities. The polynomial fit has the following parameters: $T_{\text{eff}} = 11.5 + 0.87 a_0 - 0.0035 a_0^2$. The figure on the right hand side shows the electron spectra for three different intensities after 1.41 mm propagation.

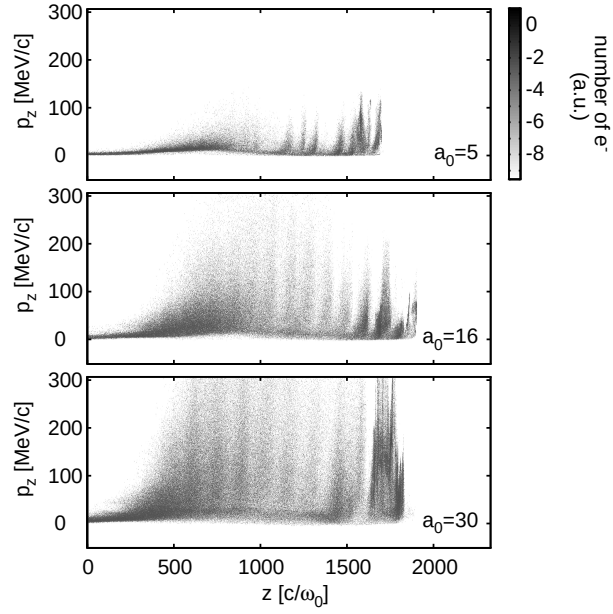


Figure 4.27: Phase space for three different laser intensities, after the laser has exited the plasma of 2 mm.

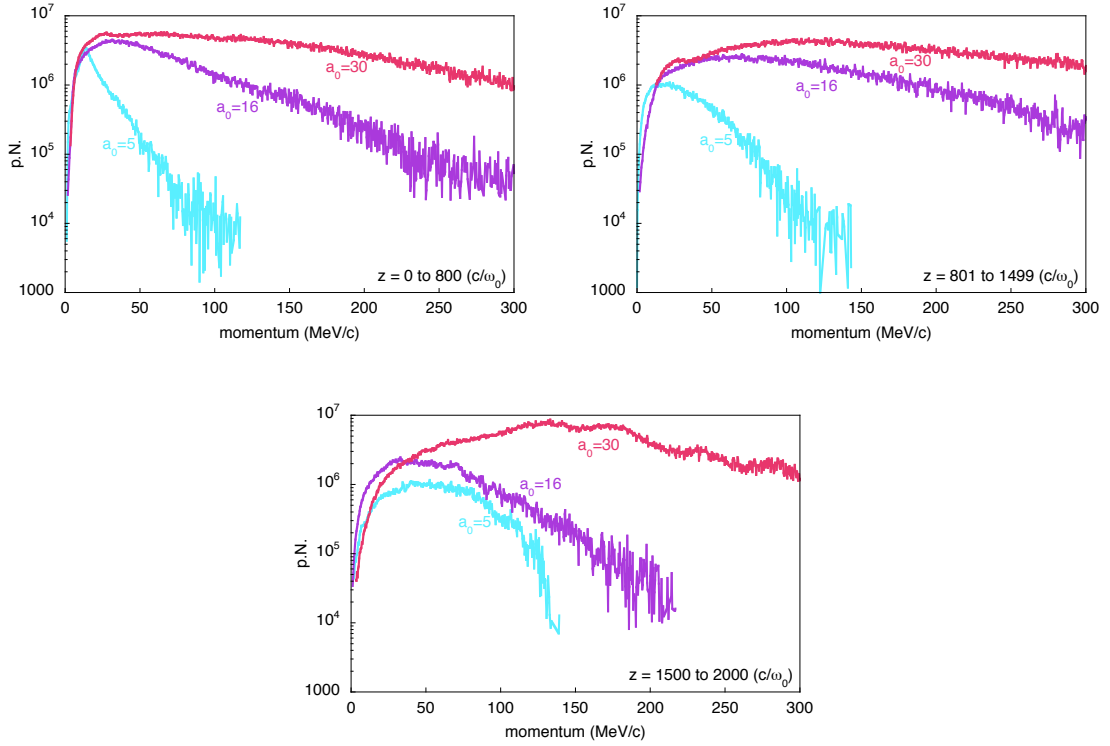


Figure 4.28: Electron spectra after 2 mm propagation for different parts of the simulation box and different intensities.

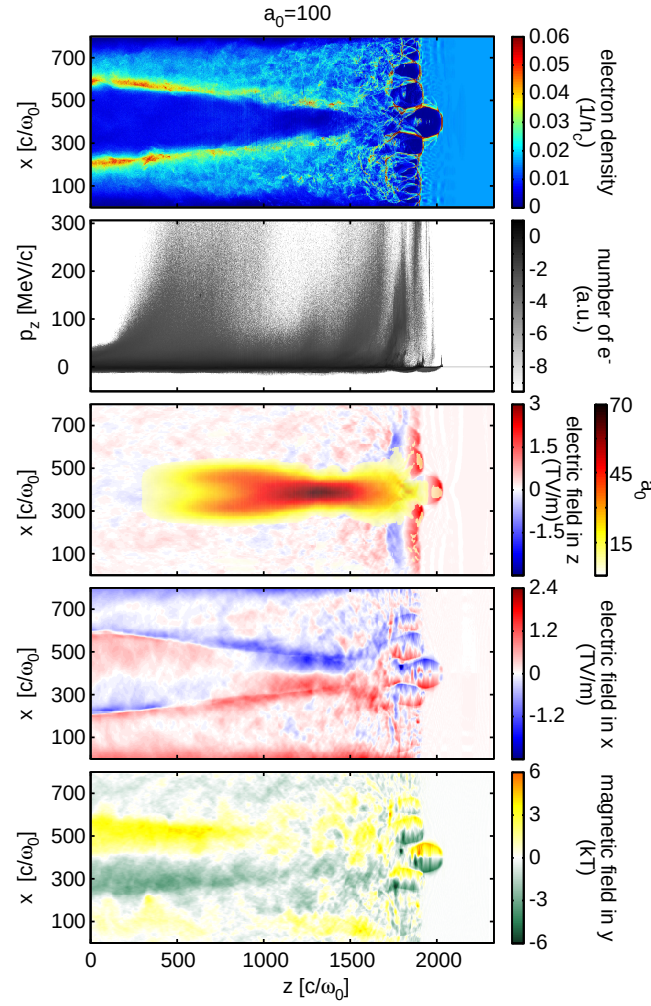


Figure 4.29: Plots of the density profile, phase space, longitudinal electric field, laser envelope, transverse electric field and perpendicular magnetic field for $a_0 = 100$, after 1.41 mm of propagation in helium with $n_e = 0.016n_c$.

Figure 4.29 shows the density profile, phase space, longitudinal and transverse electric fields and perpendicular magnetic field for $a_0 = 100$, after 1.41 mm of propagation in helium with $n_e = 0.016n_c$. $a_0 = 100$ corresponds to an intensity of $> 1 \times 10^{22} \text{ Wcm}^{-2}$, which is an order of magnitude higher than the intensities achievable with VULCAN today. SMLWFA is strongly suppressed, but stronger radial electric, and azimuthal magnetic fields lead to better, more efficient DLA. Under these condition the simulations suggest very high temperatures and electron energies in excess of 300 MeV.

4.3.2 f-number comparison

Most of the simulation parameters for the f-number comparison scan were kept the same as for the intensity scan described in section 4.3.1. The simulations were conducted for comparison with the experiment described in section 4.1.3. Here, the FWHM of the focal spots were $5 \mu\text{m}$ for f/3-focusing and $9.5 \mu\text{m}$ for the f/5 focusing geometry, leading to a_0 of 21 and 11 respectively. The electron density of the target was changed and had the following values: $7.4 \times 10^{18} \text{ cm}^{-3}$, $1.3 \times 10^{19} \text{ cm}^{-3}$ and $2.4 \times 10^{19} \text{ cm}^{-3}$, which correspond to $0.0074n_c$, $0.013n_c$ and $0.024n_c$, respectively.

The results are plotted in figures 4.30, 4.31 and 4.32, which show the density profiles, laser envelope and phase space, respectively. They are plotted so that the two columns correspond to the two different f-numbers, f/3 and f/5, and the different rows correspond to different densities, $0.0074n_c$, $0.013n_c$ and $0.024n_c$. The interaction front moves backwards (towards the left) in the box for higher densities. The simulation box is moving at c , and at the start of the simulations, the pulse is at the same position. As the simulation progresses, the pulse moves towards the left of the simulation box, with a rate that is dependent on the group velocity, which is lower for higher density[†].

Looking at the density profiles for f/3 and f/5-focusing, see figure 4.30, filamentation can be seen. For f/3-focusing, filamentation is observed at a lower density of $1.3 \times 10^{19} \text{ cm}^{-3}$ and $2.4 \times 10^{19} \text{ cm}^{-3}$, whereas for f/5 filamentation in the first 2.16 mm can only be discerned for $n_e = 2.4 \times 10^{19} \text{ cm}^{-3}$. This can be established further in the plots of the laser envelope, see figure 4.31.

At the lowest density no injection into the bubble in the first 2.16 mm is seen for both f-numbers. However, for the case of f/3-focusing, a bubble is formed whereas for the case of f/5-focusing, only the early stage of the bubble formation can be observed. For the f/3-focusing at higher densities, injected electrons can easily be seen in figure 4.30, as well as in the phase space diagrams in figure 4.32. A split of the bubble into two channels, is due to filamentation, as described above.

[†] $v_g(0.0074n_c) \approx 0.996 * c$ and for $v_g(0.024n_c) \approx 0.988 * c$.

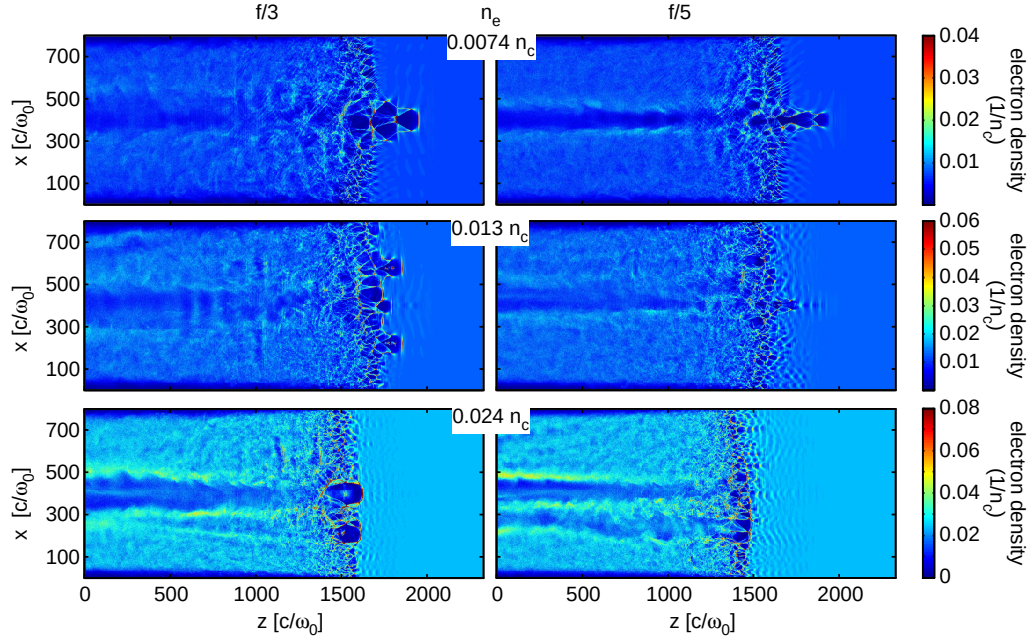


Figure 4.30: Plots of the density profile, for the two different f-numbers and three different densities for a propagation length of 2.16 mm.

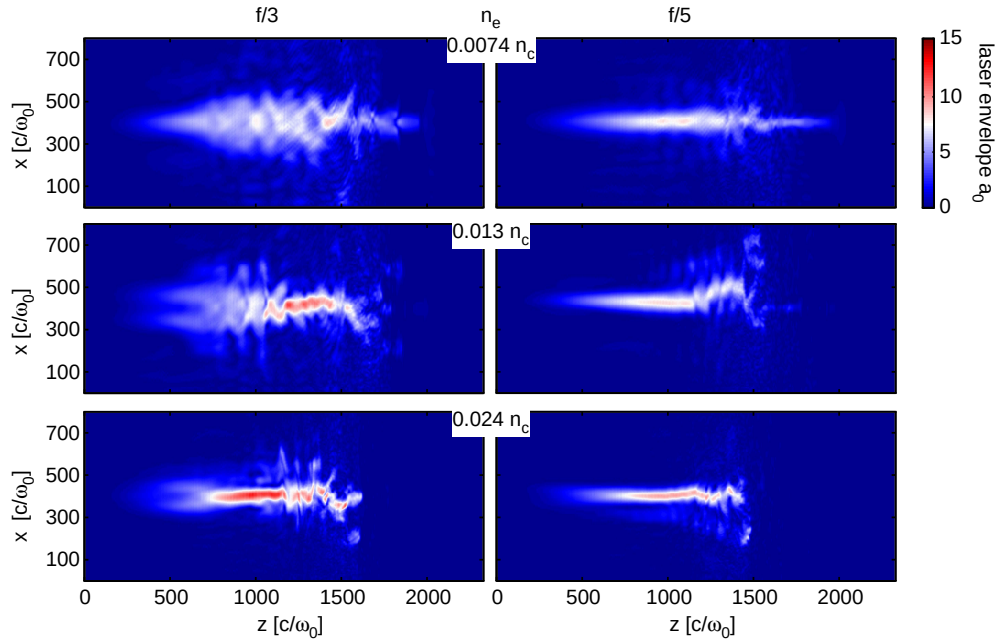


Figure 4.31: Plots of the laser envelope, for the two different f-numbers and three different densities for a propagation length of 2.16 mm.

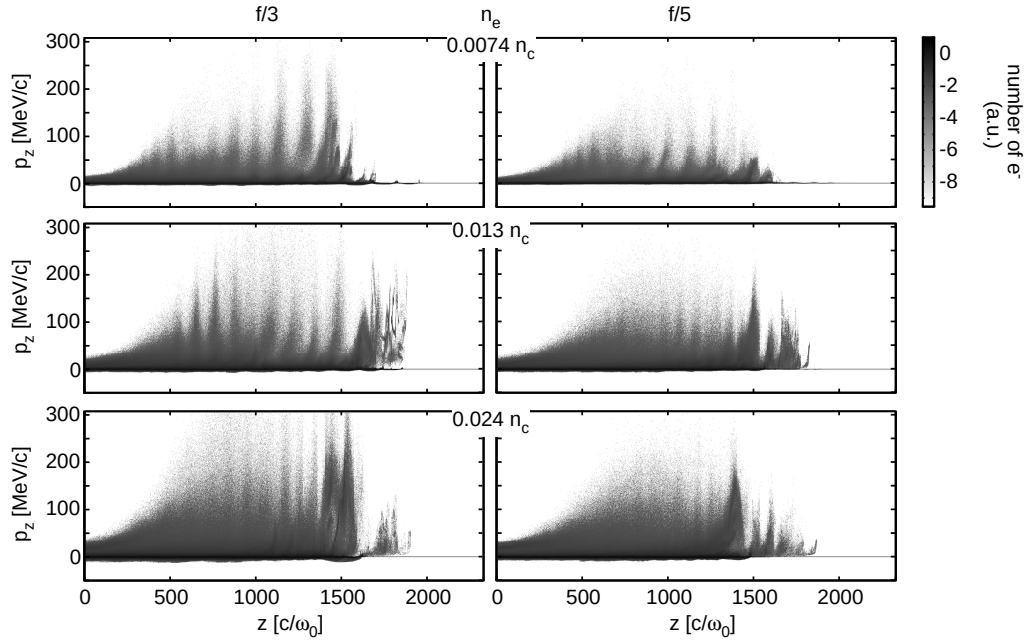


Figure 4.32: Plots of the longitudinal momentum against the position in z , for the two different f -numbers and three different densities for a propagation length of 2.16 mm.

Concerning the channel formation, it seems that for $f/5$ -focusing at low densities, the channels in figure 4.30 are more distinct.

In contrast to the experimental data, at the given propagation length of 2.16 mm, the electrons for the $f/3$ -focusing are hotter (see figure 4.33), and extend to higher energies. However, the laser was not propagated for the whole 5 mm as in the experiment, due to computational constraints, and a bigger box size would be necessary for the high densities to prevent the laser from drifting out of the box for such long interaction lengths. But a change in the electron injection into the bubble for different f -numbers was observed. By changing the f -number only slightly, the guiding of the laser was improved, as can be seen in figure 4.31.

Figure 4.33 shows the effective temperatures for different f -numbers and densities in the simulations. The temperatures are a lot higher than in the experiment, and do not show the same density and f -number dependence. This could be due to the following differences between the simulation and experiment. The simulation is mainly two dimensional, and three dimensional effects, such as the variation of

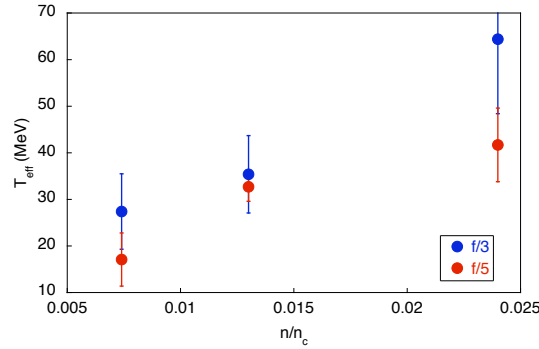


Figure 4.33: Temperatures for different f -numbers and densities from the simulations.

the intensity with spot size is therefore different. In particular for a diverging beam as in the $f/3$ case, the a_0 is overestimated as the spot size gets bigger, because it expands only in one transverse dimension, not two. Furthermore, the focal spot in the simulation is an optimal, smooth spot, whereas in the experiment hot spots in the focus can appear. This would stimulate filamentation further which is detrimental to the acceleration. In the simulations the focal spot size was set to be the same as in the experiment, since the simulation assumes a diffraction limited spot. Hence the Rayleigh length is longer in the simulation. The difference between f -numbers, can also result in a change in the focal spot quality. In the experiment the focal spot quality improves for larger f -number, whereas in the simulation, because of the smooth focal spots, the difference in laser intensity is less important.

4.3.3 Length scan

The simulations in this section were conducted as a comparison with the interaction length scan using the f/10-focusing, described in section 4.2. The simulation box is co-propagating with the laser pulse at the speed of light, c . A 500 fs long pulse with a gaussian temporal shape, and linear polarisation in the plane of the simulation box, was focused to the edge of fully ionised helium plasma with a density of $8 \times 10^{18} \text{ cm}^{-3}$. The simulations were run for a focal spot size of FWHM $32 \mu\text{m}$, and the laser had an $a_0 = 1.1$. The resolution of the $783 \mu\text{m}$ by $134 \mu\text{m}$ big simulation box was 118 nm ($\approx \lambda_p/100$) in the z - or propagation-direction and 200 nm ($\approx \lambda_p/59$) in the transverse direction. In this simulation, the box size was chosen to be able to accommodate the laser pulse as it moves back in the box for longer propagation lengths, due to the mismatch of group velocity and c . The time interval the data was saved at was 50 fs, and the transverse boundary conditions were open space.

Figure 4.34 summarises the density profile, phase space and laser envelope at three different propagation lengths. Note that the parts of the simulation box plotted here are chosen to contain the laser pulse. They are all $2330 (c/\omega_0)$ long, which is half the size of the complete simulation box, and start at the following positions of the full one: For 3 mm the plot starts at $2000 (c/\omega_0)$, for 5 mm the box starts at $1890 (c/\omega_0)$ and the box for the case of 8 mm propagation begins at $1500 (c/\omega_0)$.

For the propagation of 3 mm, it is possible to observe the modulation of the laser envelope at the plasma frequency. For longer interaction lengths, the depletion of the laser intensity becomes obvious. Furthermore, as the laser propagates the plasma wave structure in the density profiles becomes less pronounced, and the dephasing of the electrons is apparent in the phase space diagrams.

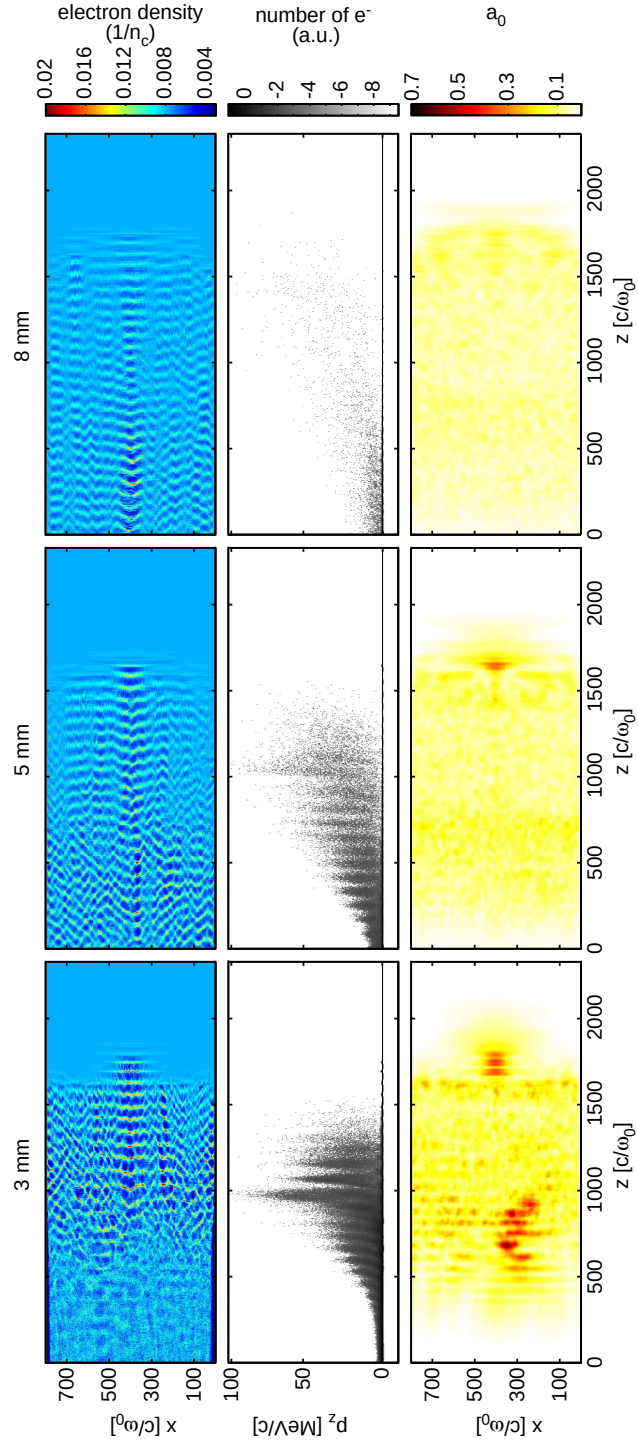


Figure 4.34: Plots of the density profile, phase space and laser envelope for three different propagation lengths in helium with $n_e = 0.008 n_c$.

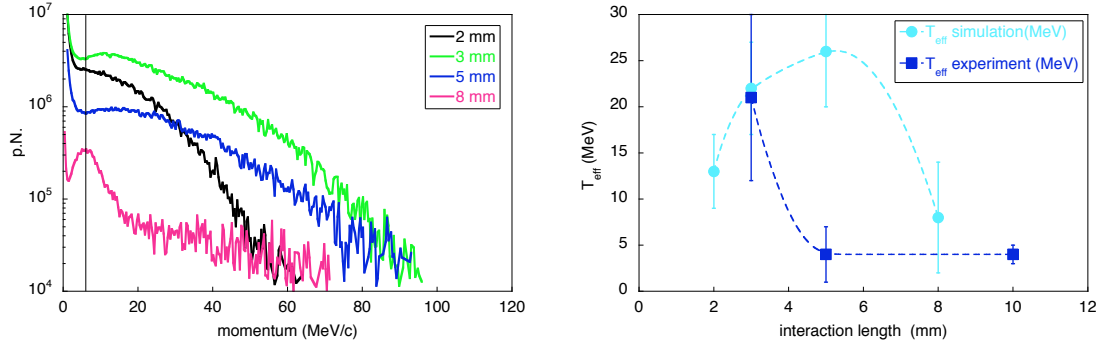


Figure 4.35: Electron spectra from the simulations at different propagation lengths. The line at 6 MeV indicates the detector limit in the experiment. The plot on the right hand side shows the effective temperatures from the experiment and the simulations plotted against interaction length.

Electron spectra for different propagation lengths in the simulation are plotted in figure 4.35. The vertical line indicates the low energy detector limit in the experiment described in section 4.2. At a propagation length of 3 mm, the electron spectrum has a fairly long ridge and extends to about 90 MeV, similar to the one in the experiment, see figure 4.20. For longer interaction lengths, the number of electrons decreases, indicating dephasing. This is also in agreement with the experimental measurements. Though in comparison to the experiment, initial dephasing at 5 mm is less, but increases rapidly for 8 mm.

Chapter 5

Electrons from ultra-thin targets

The goal of the experiment described in this chapter was to study the interaction of thin foil targets at intensities greater than 10^{20} Wcm^{-2} . To do this, plasma mirrors were required, to stop the target being destroyed prior to the arrival of the main pulse. These measurements are relevant to the radiation pressure acceleration (RPA) scheme of ion acceleration, where the whole bulk is accelerated if the target is thin enough.

5.1 Experimental setup

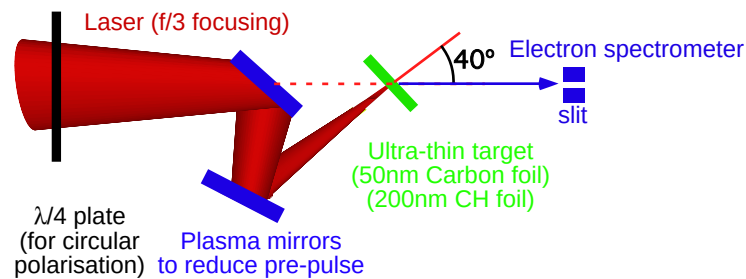


Figure 5.1: Sketch of the experimental setup for high contrast interactions with ultra-thin targets.

The experiment described, was performed using the VULCAN Petawatt laser. For the shots discussed here, the laser had a duration of (520 ± 100) fs and an energy of (60 ± 7) J on target. This “low” energy was due to the use of double plasma mirrors,

which lead to an energy transmission of around 10%. An $f/3$ off-axis parabolic mirror focused the laser via a pair of plasma mirrors onto ultra-thin targets (50 nm carbon foils and 200 nm thick plastic foils). Focal spot measurements at $3\omega_0$ gave a FWHM of approximately $8\ \mu\text{m}$ [162]. This implies intensities of about $1.5 \times 10^{20}\ \text{Wcm}^{-2}$, and $a_{0\text{lin}} = 11$ for linear polarisation, where a_0 is the peak normalised vector potential. To produce circular polarisation, a $\lambda/4$ -plate was put in the focusing beam, since the collimated laser beam diameter was larger than the $\lambda/4$ -plate. The peak a_0 for circular polarisation can be inferred to be $a_{0\text{circ}} = 7.8$, since $a_{0\text{circ}} = a_{0\text{lin}}/\sqrt{2}$.

Figure 5.1 shows a sketch of the experimental setup, where plasma mirrors are plotted in blue. The targets were shot at normal incidence, and the spectrum of energetic electrons produced in the interaction was measured 40° from the laser propagation direction. This angle is a result of the plasma mirror setup, since the electron spectrometer was initially set up along the laser propagation direction without the plasma mirrors. A 2 mm slit was placed in front of the entrance to the electron spectrometer, to limit the angular acceptance and hence improve the energy resolution. In this experiment the electron spectrometer was run with a field strength to measure electron energies between 1.2 and 30 MeV.

5.2 Experimental data

Figure 5.2 shows the experimentally obtained electron spectra. The number of electrons per MeV per steradian is plotted against the electron energy in MeV. The light blue and violet data correspond to linear laser polarisation, and the dark blue and red ones to circular polarisation. For both target thicknesses, the maximum energy for the linear case is about twice as high as for the circular case, and for the spectra for linear polarisation almost twice the effective temperature (T_{eff}) is observed.

This can be seen in figure 5.3, which shows the effective temperatures against target thickness, for both laser polarisations. The temperatures from the “manual” fit give the higher electron energies the same weight as the lower energies. Whereas the usual fit favours the temperature of the lower energies, because there are more

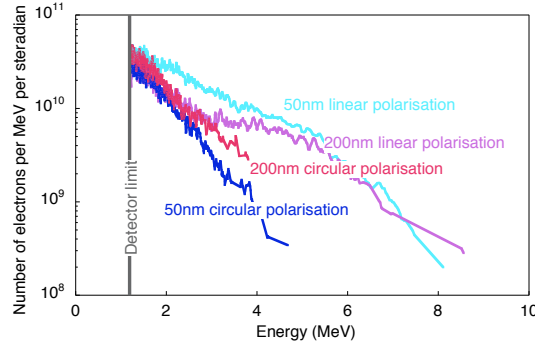


Figure 5.2: Electron spectra for ultra-thin targets, taken at 40° off the laser propagation.

data points to take into account. This is due to the way the electron spectrometer deflects the electrons and its spatial resolution in the deflection direction. Nevertheless they are in good agreement, and both show the trend discussed in the preceding paragraph. The two different thicknesses show very similar behaviour, and the temperatures are comparable within their error margin. This indicates, that the different thicknesses of the targets did not change the interaction regime significantly.

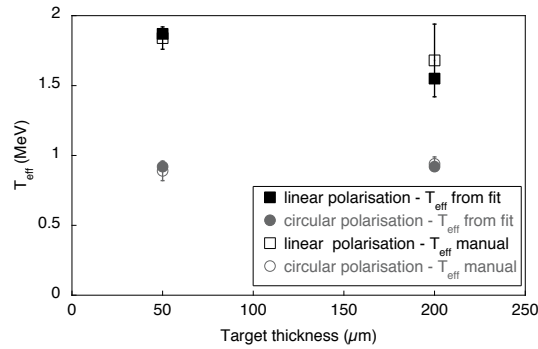


Figure 5.3: Temperatures of the electron spectra in figure 5.2.

Comparing the effective temperatures to the ponderomotive potential ($m_e c^2 (\langle \gamma \rangle - 1) = 3.5 \text{ MeV}$)*, shows that the electrons measured off axis feature a lower T_{eff} than

*where $\langle \gamma \rangle = \sqrt{1 + \frac{1}{2} a_{0\text{lin}}^2}$ for linear polarisation and $\langle \gamma \rangle = \sqrt{1 + a_{0\text{circ}}^2}$ for circular polarisation.

theoretically predicted for the forward direction. Furthermore the predicted maximum momentum $m_e c \gamma = m_e c (1 + a^2/2) \approx m_e c (1 + a_0^2/4) \approx 16 \text{ MeV}/c$, for linear polarisation in the forward direction also exceeds the maximum energy as measured at 40° off laser axis.

5.3 Simulations on the polarisation comparison for ultra-thin targets

In order to help gain insight into the interaction, 2D simulations, using the particle-in-cell code OSIRIS [165–167], were performed (see section 4.3). The simulations described here were conducted to compare to the measurements from ultra-thin targets described in the previous section. A stationary box, with a target of 200 nm thickness, situated $1.68 \mu\text{m}$ from the left boundary, was simulated. The simulation box had dimensions of $13.4 \mu\text{m}$ by $13.4 \mu\text{m}$, with resolutions of 1 nm ($\approx \lambda_p/31.4$) in the laser propagation direction and 5 nm ($\approx \lambda_p/6.3$) in the perpendicular direction. The laser propagated in the box from the left boundary to the right, and was polarised in the plane, for linear polarisation. Data was saved for every 10 fs interval, and the simulation was run with 4 particles per cell for both particle species (electrons and ions). The target electron density was set to correspond to $1000n_c$ and the ion species was carbon⁶⁺. The focal spot FWHM was $8 \mu\text{m}$, and a_0 was set to 10, for linear polarisation. For circular polarisation the pulse in the simulations was set up by two linearly polarised waves, one of them was dephased by $\lambda/4 \hat{=} \pi/2 = 90^\circ$, and had a polarisation out of the plane of observation. The value of a_0 for the two pulses in the case of circular polarisation was: $a_{0\text{ circ}} = a_{0\text{ lin}}/\sqrt{2}$.[†] The pulse shape was gaussian in time with a FWHM of 500 fs.

[†]The energy in a pulse is $\mathcal{E} \propto a_0^2 \tau \omega_0^2$, since the circularly polarised pulse can be described as two linearly polarised pulses: $\mathcal{E} = \mathcal{E}_1 + \mathcal{E}_2 \Rightarrow \mathcal{E} = (a_{0,1}^2 + a_{0,2}^2) \tau \omega_0^2 \rightarrow a_0^2 \tau \omega_0^2 = (a_{0,1}^2 + a_{0,2}^2) \tau \omega_0^2$ for circular polarisation $a_{0,1} = a_{0,2} = a_{0\text{ circ}}$ and one obtains: $2a_{0\text{ circ}}^2 = a_{0\text{ lin}}^2 \Rightarrow a_{0\text{ circ}} = a_{0\text{ lin}}/\sqrt{2}$

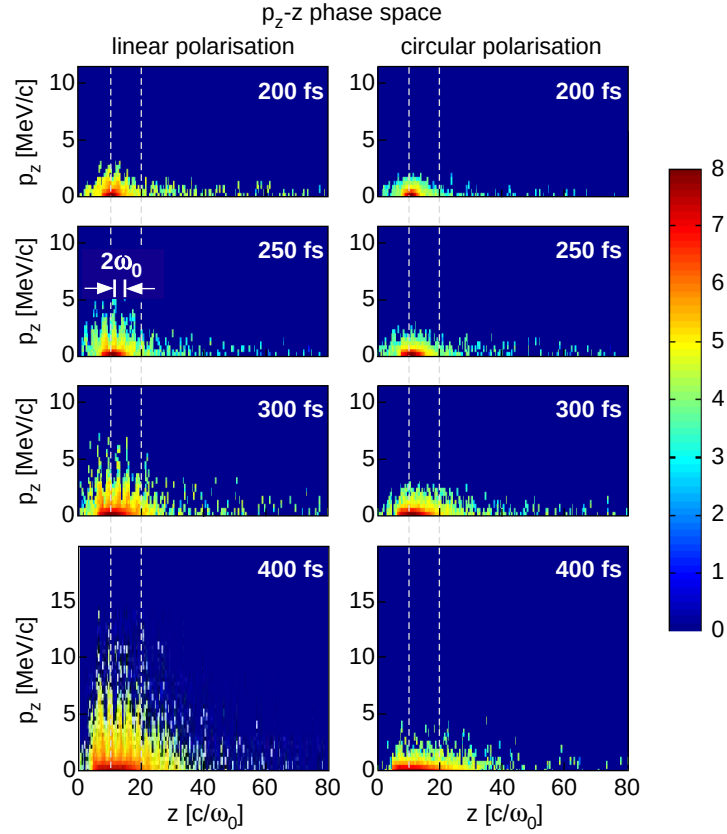


Figure 5.4: $p_z - z$ phase space for different times in the simulations, the numbers in the upper corners indicate the time in the simulation. The left column corresponds to a linearly polarised pulse, whereas the right column shows the phase space for circular polarisation. Linear polarisation shows bunching of the electrons at $2\omega_0$. For circular polarisation the $2\omega_0$ bunching is not present. Note: the original target starts at $10 c/\omega_0$. And $c/\omega_0 = 0.168 \mu\text{m}$. The colourbar indicates the number of electrons in arbitrary units.

Figure 5.4 shows the $p_z - z$ phase space for different times (rows) and different laser polarisations (columns). The numbers in the top right corners of the images indicate the time in the simulation in fs. The dashed grey vertical lines correspond to the positions of 10 and 20 (c/ω_0) in the simulation box, the target front surface is also situated at position 10 (c/ω_0) = 1.68 μm . For the case of linear polarisation the heating of the electrons due to $\mathbf{j} \times \mathbf{B}$ heating can be observed, since the electrons are clearly bunched at $2\omega_0$. This effect is suppressed for the case of circular polarisation, as described in section 2.3.1.4, leading to a decrease in the electron temperature.

The maximum electron momenta achieved in the simulation in the forward direction, are ≈ 16 MeV for linear polarisation and ≈ 5 MeV for circular polarisation. These are very close to the theoretically predicted momenta $m_e c \gamma \approx 16$ MeV/c for linear polarisation and ≈ 7 for circular polarisation.

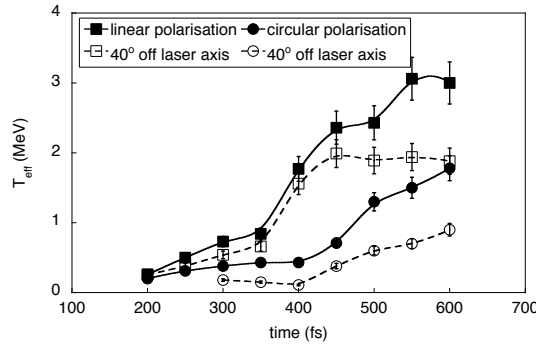


Figure 5.5: Temporal evolution of the effective temperature in the simulations for linear and circular polarisation in the forward direction, as well as for electrons emitted at 40° off laser axis, for a 200 nm thick foil.

In figure 5.5 the temporal evolution of the effective temperature is plotted for different scenarios. Linear and circular polarisation can be compared, as well as the whole electron population in the forward direction and only electrons that are traveling with an angle of 40° off laser axis. Circular polarisation leads to a lower temperature than linear polarisation. It can be seen that, for both linear and circular polarisation, the effective temperature at 40° is about a factor of 1.5 smaller than the ones on axis. The temperatures at 40° are close to the ones measured experimentally, and the ones in the forward direction for linear polarisation approach

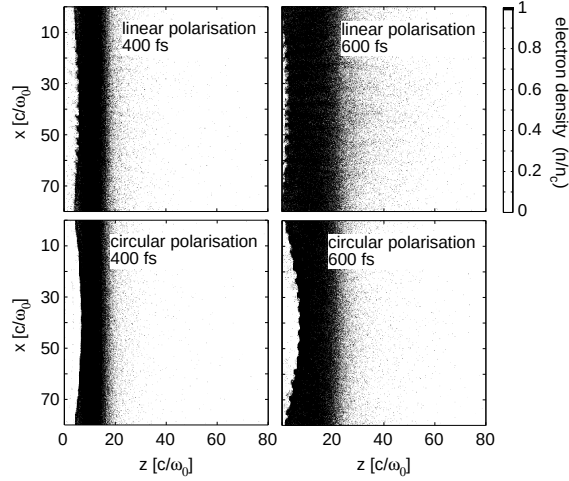


Figure 5.6: Density profiles showing the density above critical for two different times in the simulations and linear and circular polarisation.

the ponderomotive potential, $U_p = 3.1$ MeV.

Looking at the density profiles for the two different polarisations at two different times in the simulation, shown in figure 5.6, one can observe that for circular polarisation, the foil is displaced more uniformly, due to the constant ponderomotive push. The initial electron density in the simulation ($1000n_c$), was approximately the electron density in the real carbon foil. From the temperature measurements in figure 5.5 we can observe that the heating for circular polarisation increases after 400 fs.

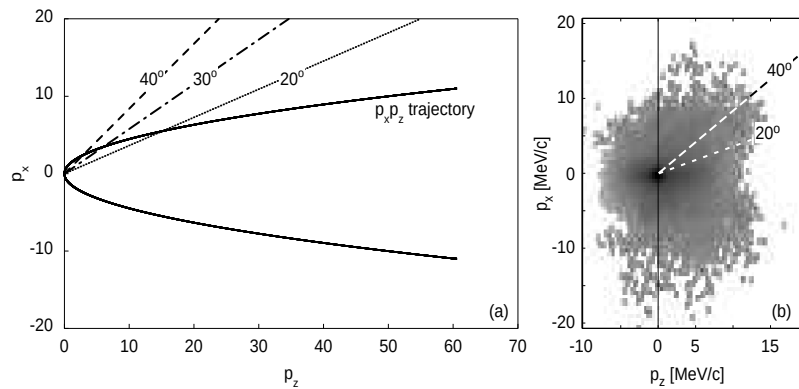


Figure 5.7: Electron trajectory in momentum space, theoretical prediction (a), and from simulation (b).

Figure 5.7(a) shows the trajectory in momentum space for an electron in a linearly polarised electromagnetic field, for the given $a_0 = 11$ (see appendix A.1). Here z is the laser propagation direction and x is the direction of laser polarisation. The lines indicate different deflection angles away from the laser axis. In figure 5.7(b), the electron momentum space from the simulation for linear polarisation, at 400 fs into the simulations is plotted. The parabolic trajectory is filled with electrons in the simulation, which is due to the ponderomotive push, that also accelerates electrons in the case of circular polarisation.

For linear polarisation a simple relation between the Lorentz factor of the electron and the angle θ between the electron trajectory and the laser propagation axis can be deduced from equations (A.16)[‡] and $(\gamma - 1)m_e c^2 = cp_z(t)$.

$$\tan \theta = \sqrt{\frac{2}{\gamma - 1}},$$

which for an angle of 40° gives a γ of 3.8 leading to a ponderomotive potential of around 1.1 ± 0.2 MeV. This underestimates the effective temperature observed in the experiment and simulation for linear polarisation.

Looking at the ion spectra plotted in figure 5.8, for (a) linear and (b) circular polarisation, shows that the shock acceleration is more pronounced for circular polarisation. This can be seen by the strongly peaked spectra for circular polarisation at times where the ponderomotive potential of the laser is highest. The ion energies these “kinks” correspond to are 20 MeV at 600 fs, 36 MeV at 700 fs and 56 MeV at 800 fs.

The low intensity part of the laser accelerates electrons, as they leave the target, they set up an electric sheath field, which accelerates ions (TNSA). When the laser reaches a higher intensity, shock acceleration of the ions can be observed, but only for circular polarisation, because of the constant ponderomotive push. Linear polarisation features a larger divergence, compared to circular polarisation. This can be seen in figure 5.9, where the ion phase and momentum space are plotted for linear and circular polarisation at 600 fs.

[‡] $p_z = \frac{p_x^2}{2m_e c}$

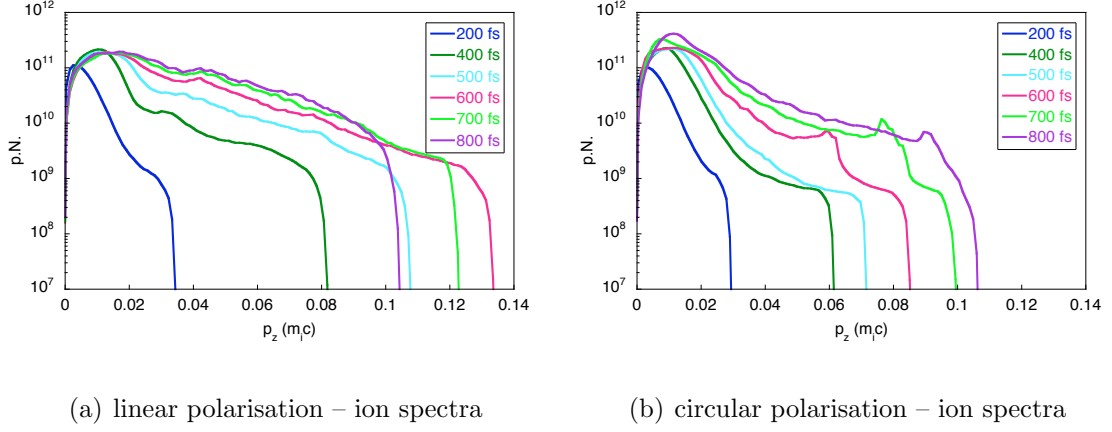


Figure 5.8: The ion spectra, for ions moving in the laser direction, from the simulations of a 200 nm thick target and for different time steps in the simulation. (a) for linear polarisation and (b) for circular polarisation. The decrease in maximum ion momentum for later times in the case of linear polarisation is due to ions leaving the simulation box.

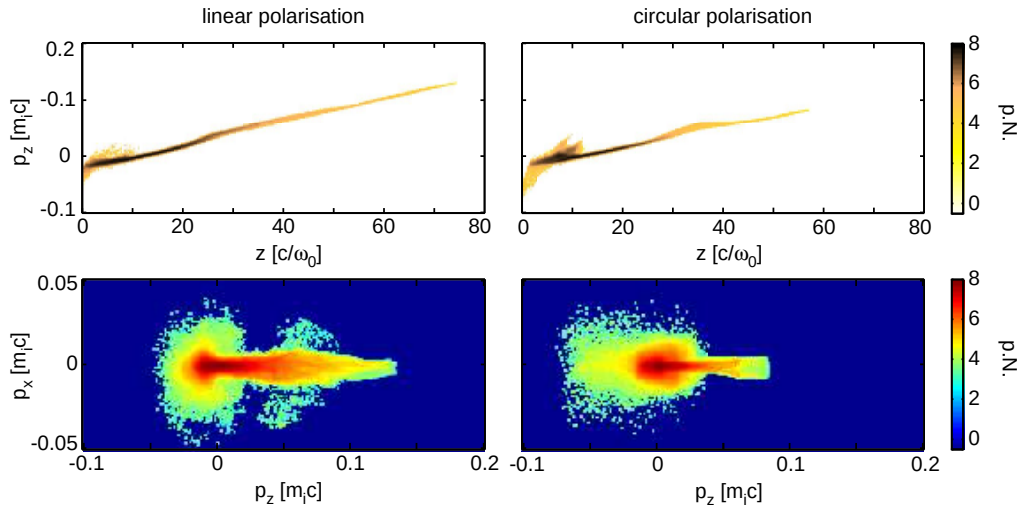


Figure 5.9: Ion phase and momentum space at 600 fs into the simulation.

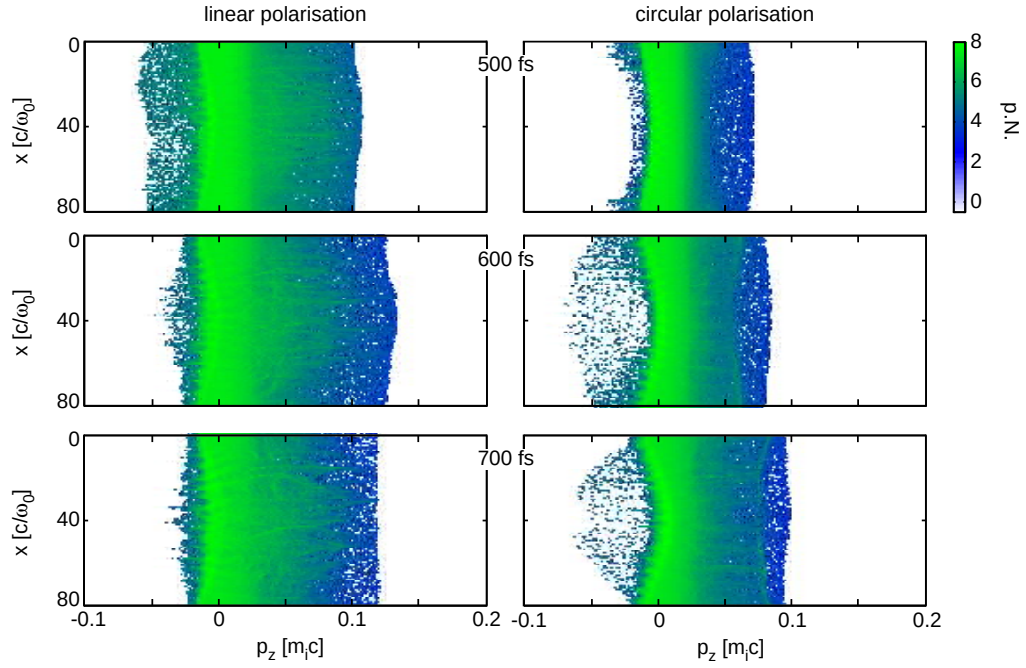


Figure 5.10: Ion transverse position plotted against longitudinal momentum for linear and circular polarisation (columns) at three different times in the simulations (rows), namely 500, 600 and 700 fs. For the circular polarisation at 600 and 700 fs the “shock” accelerated ion front can be observed.

Figure 5.10 shows the transverse position of the ions plotted against the longitudinal momentum for linear and circular polarisation at three different times in the simulation. For circular polarisation at later times, 600 and 700 fs, the ion “shock” front leading to the kink in the spectrum can be observed.

Chapter 6

Electron acceleration from “thick” targets

In this chapter electron acceleration from “thick” overdense plasmas is investigated, and it is organised in the following way. First we study the effect of different contrast ratios for 20 μm thick copper foils. Then a scan is performed using Au foil targets ranging from thicknesses of 10 to 100 μm . Both linear and circular polarisation are investigated over this target thickness range. The different experimental findings are compared in section 6.4. This is followed by simulations studying the effects of different laser polarisation and varying target thickness.

6.1 Contrast ratio comparison

A sketch of the experimental setup can be seen in figure 6.1. The single plasma mirror, which improves the contrast ratio, could be exchanged for a highly reflective (HR) mirror, to keep the lower contrast ratio. Here the energy on target was (295 ± 13) J, leading to an intensity of $8.3 \times 10^{20} \text{ Wcm}^{-2}$ ($a_0 = 26$, $U_p = 8.9 \text{ MeV}$), and the pulse duration was (700 ± 50) fs. The targets were flat $20 \mu\text{m}$ thick copper foils, which were shot at normal incidence.

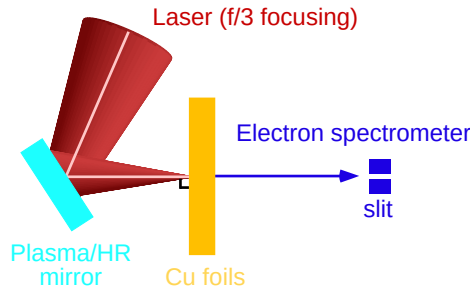


Figure 6.1: Sketch of the experimental setup for the contrast ratio comparison.

Figure 6.2 shows the backlit side image of the interaction at (10 ± 5) ps before the main pulse arrives. The laser is coming in from the right. For the low contrast ratio 6.2(a) a preplasma can be observed, which has a scale length of $5 \mu\text{m}$ [169]. The bright spot in figure 6.2(b) is due to self-emission during the interaction. But the backlit image shows that there is hardly any pre-plasma for the case of high contrast ratio.

Figure 6.3 shows the corresponding electron spectra for shots comparing a low contrast ratio, $\approx 10^{-7}$ (red curve), with a high contrast ratio, $\approx 10^{-9} - 10^{-10}$ (blue curve). The temperatures for the spectra are 7.9 and 4.3 MeV for low and high contrast, respectively. Both temperatures are lower than the ponderomotive potential, though barely in the case with large prepulse. As expected for a low contrast ratio (due to the formation of a pre-plasma) the observed number of electrons, maximum energy and effective temperature are higher than for high contrast ratio. This can be explained by underdense acceleration mechanisms starting to play a role for the case of low contrast ratio. Additionally self-focusing in the underdense plasma

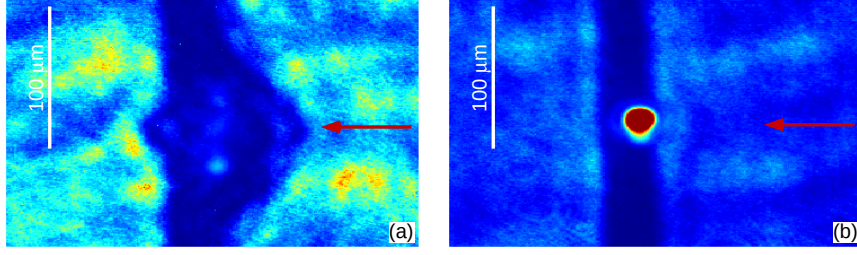


Figure 6.2: Probe images for two different contrast ratios. (a) using the HR mirror, leading to a low contrast, and (b) using the plasma mirror, leading to a high contrast. The red arrow indicates the laser direction. (Figure courtesy of J. Schreiber)

could lead to an increase in laser intensity. Whereas for the high contrast ratio a steeper density ramp is expected and overdense acceleration mechanisms are more important.

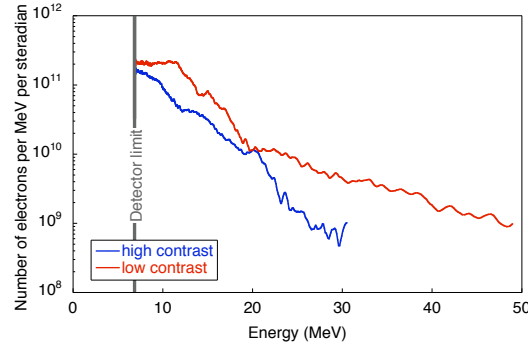


Figure 6.3: Electron spectra for different contrast ratios, using 20 μm thick copper foils at normal incidence.

The data in the following sections was taken with the lower contrast ratio, therefore with a pedestal intensity of $\approx 1 \times 10^{14} \text{ Wcm}^{-2}$.

6.2 Normal incidence irradiation of gold foils

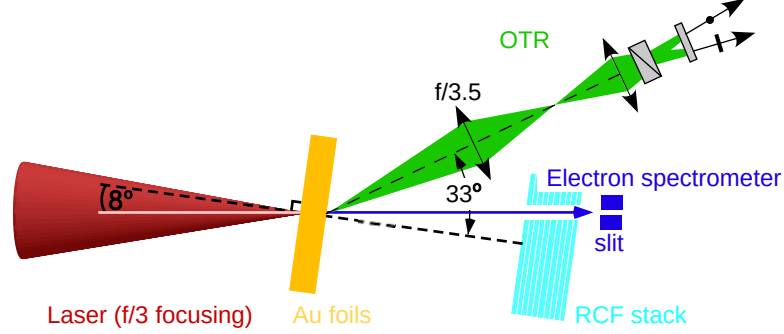


Figure 6.4: Sketch of the experimental setup for normal incidence irradiation of gold foils.

Figure 6.4 shows a sketch of the experimental setup. For this experiment the OPCPA front end was used to reduce the pulse length in TAW. Therefore the laser had a duration of (642 ± 124) fs indicating a pulse length of $c\tau = (193 \pm 37)$ μm and an energy of (55 ± 4) J on target. An $f/3$ off-axis parabolic mirror focused the laser onto the solid targets. Spot size measurements gave a FWHM of 9.5 μm . This implies intensities of $(8.3 \pm 2.2) \times 10^{19}$ Wcm^{-2} . Therefore the amplitude of the normalised vector potential, a_0 , for linear polarisation was (8.2 ± 1.1) . The laser polarisation could be changed from linear to circular, by putting a $\lambda/4$ wave plate into the collimated beam. The angle of incidence was 8° off target normal. This was a compromise between trying to shoot along target normal and preventing back reflections into the laser system. The electron spectra were measured along the laser axis using a magnetic spectrometer, and imaging plates were used as the electron detector. Here the targets were gold foils with thicknesses between 10 and 100 μm .

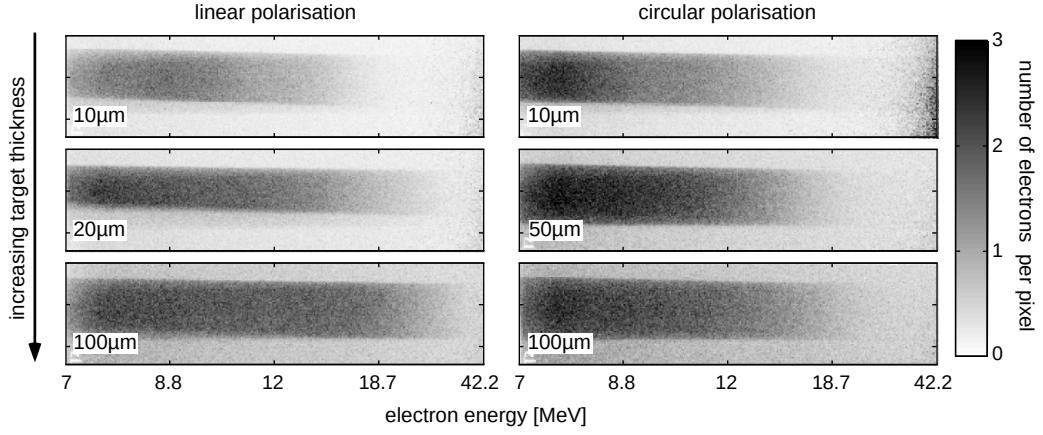
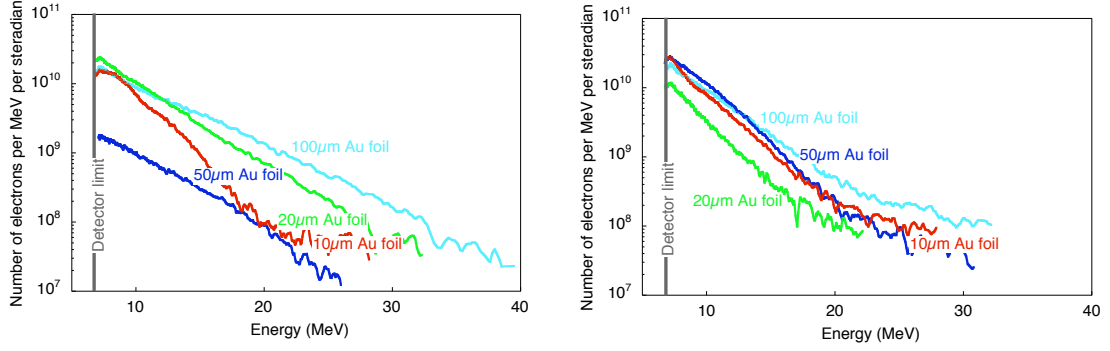


Figure 6.5: The imaging plate data. The different signal widths are due to the radiochromic film (RCF) stack, pictured in figure 6.4, clipping some of the detected electron beam.

The raw imaging plate data is shown in figure 6.5 and the corresponding electron spectra are shown in Figure 6.6, (a) linear laser polarisation and (b) circular laser polarisation. Here, the number of electrons per MeV per steradian is plotted against the electron energy in MeV. The different colours correspond to different target thicknesses. As can be seen in figure 6.5, the electron beams look very smooth, in comparison to the quasi mono-energetic electron spectra in the LWFA regime [15–17]. In general thicker targets show a higher number of electrons and charge, see figure 6.7 (exceptions: 50 μm linear and 20 μm circular). This trend is even more pronounced for higher electron energies. For the two different polarisations, however, the total charge in the spectra is comparable.

The shots for the 50 μm thick target for linear polarisation and 20 μm for circular polarisation have anomalously low electron numbers. This could be due to a change in beam pointing and therefore the effective temperature is a more reliable measurement. The temperatures of the spectra were obtained by fitting a quasi maxwellian function, $y = A * \exp(-E/T_{\text{eff}})$, to the spectra.

Figure 6.8 shows the effective temperature vs. the target thickness for the two different polarisations. The effective temperatures measured are of the order of that predicted by Wilks [125] (see section 2.3.1.4). With the wavelength and the intensity in this experiment, one would expect: $T_{\text{hot}} = (\sqrt{1 + 0.74 I_{18} \lambda_{\mu}^2} - 1) \times 511 \text{ keV}$



(a) Electron spectra – linear polarisation

(b) Electron spectra – circular polarisation

Figure 6.6: Electron spectra for linearly (a) and circularly (b) polarised laser. The number of electrons per MeV per steradian is plotted versus the electron energy in MeV. Different colours correspond to different target thicknesses.

≈ 3.5 MeV. Wilks assumes that the energy associated with the electrons is given by the potential they experience during the interaction. Therefore this temperature scaling is also called ponderomotive scaling.

The grey box in figure 6.8 indicates the target thicknesses that are thicker than half the laser pulse length. In figure 6.8 two trends can be observed:

- (1) larger T_{eff} for linear polarisation.
- (2) increase of T_{eff} with increasing foil thickness.

These trends can also be seen in the raw data in figure 6.5 and are further discussed in sections 6.4.1 and 6.4.2.

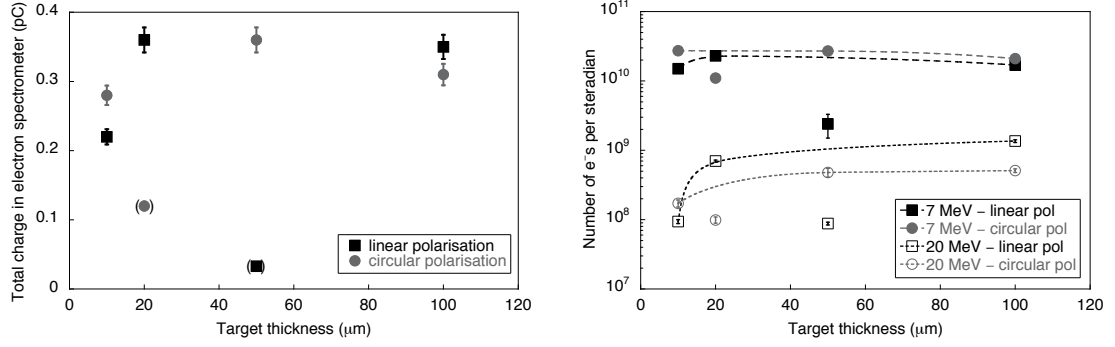


Figure 6.7: Charge in the electron spectrometer for different thicknesses and laser polarisations (left). And the number of electrons per steradian for two different electron energies and both polarisations (right).

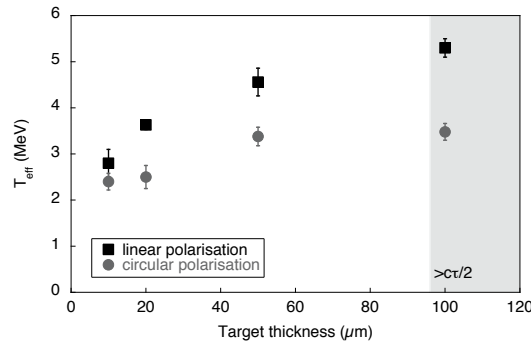


Figure 6.8: Effective temperatures plotted against target thickness for linear (■) and circular polarisation (●). The grey box indicates the thicknesses for which the target is thicker than half the laser pulse length. The error bars were obtained by a manual fit, which weighs the energies equally.

6.3 Off-target-normal incidence irradiation of gold foils with different thicknesses

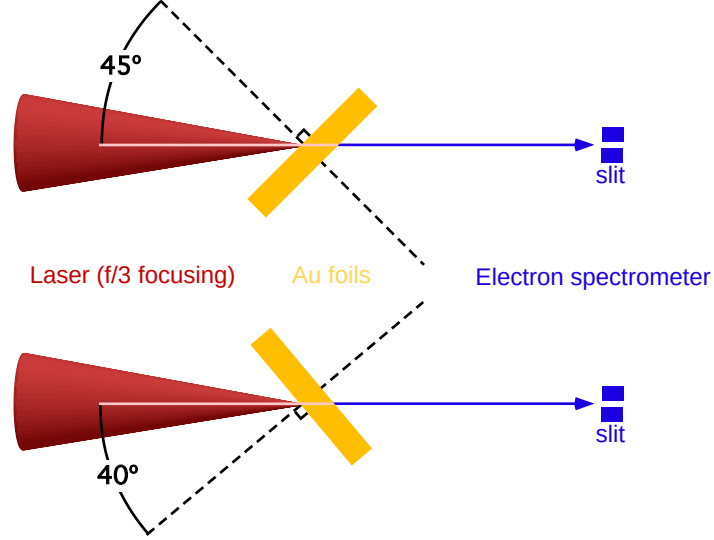
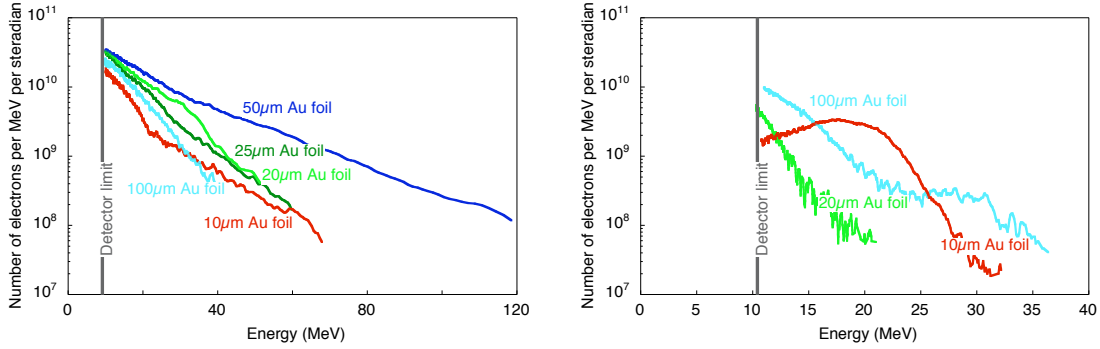


Figure 6.9: Setups for off-normal-incidence measurements of hot electrons using thick targets.

The increase of T_{eff} with target thickness was also investigated on two experimental runs with varying laser incidence angles. The laser was incident onto gold foil targets ($10 - 100 \mu\text{m}$ thick and $5 \text{ mm} \times 5 \text{ mm}$ in dimension) at 40° and 45° , as can be seen in figure 6.9, and the electron spectra were measured along the laser propagation direction. The FWHM of the focal spot were $8 \mu\text{m}$ and $5 \mu\text{m}$, and the pulse length was $(517 \pm 103) \text{ fs}$ and $(630 \pm 120) \text{ fs}$, respectively. For the shots discussed, the laser energy on target was $(299 \pm 25) \text{ J}$ for the 40° , leading to an intensity of $I \approx 7.2 \times 10^{20} \text{ Wcm}^{-2}$ ($a_0 \approx 24$). In the case of 45° the energy on target was $(208 \pm 44) \text{ J}$, corresponding to an intensity of $7.6 \times 10^{20} \text{ Wcm}^{-2}$, and $a_0=25$.

The spectra for the case of 40° incidence can be seen in figure 6.10(a). Again, a thickness dependence of the effective temperature can be observed. The shot for the $100 \mu\text{m}$ target had a lower laser energy on target ($<250 \text{ J}$, $a_0=22.4$ and $I = 6 \times 10^{20} \text{ Wcm}^{-2}$), which may explain why the number of electrons as well as the maximum electron energy are lower. For the $50 \mu\text{m}$ thick target electrons in excess of 100 MeV can be observed.

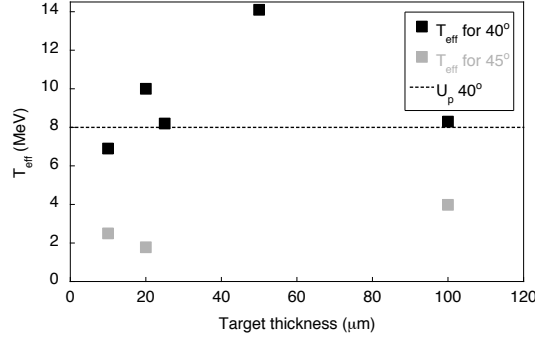
Pre-plasma measurements from this experiment show, that the critical surface moves out by $8 - 10 \mu\text{m}$, and underdense plasma can be observed up to $50 \mu\text{m}$ in front of the target, by the time the main pulse arrives [170]. The temperatures for the 40° , shown in figure 6.11, approach the ponderomotive potential. The higher temperatures could possibly be due to self-focusing of the laser, which would decrease the spot size and increase the intensity on target. Also electron acceleration due to underdense mechanisms due to the long prepulse are likely to contribute to the high energies produced.



(a) Au foil thickness scan with 40° incidence. (b) Au foil thickness scan with 45° incidence.

Figure 6.10: Electron spectra for the 40° incidence (a) and 45° incidence (b) experiments. The different colours correspond to different thicknesses, the colour scheme is the same as for figure 6.6.

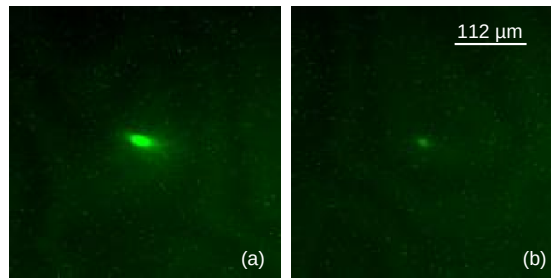
The spectra in figure 6.10(b) were taken with the 45° incidence geometry. For the $10 \mu\text{m}$ thin target in figure 6.10(b), the spectrum is non-maxwellian. For the thicker gold targets (20 and $100 \mu\text{m}$) the trend seems to be similar to the ones described in section 6.2, namely higher T_{eff} and number of electrons for increasingly thicker targets. The temperatures for the spectra are plotted in figure 6.11. The fact that the ones for 45° are not close to the ponderomotive potential, suggests that the intensity was lower than the one expected, possibly due to filamentation. This is further enhanced by the non-maxwellian shape of the spectrum for the $10 \mu\text{m}$ thick target, which suggests a large underdense region.

Figure 6.11: Plot of T_{eff} for off axis incidence experiments.

6.4 Summary of experimental results

6.4.1 Linear and circular polarisation comparison

The dependence of the effective temperature on laser polarisation has been seen in two experiments, described in chapters 5 and 6.2. A decrease of maximum electron energy as well as effective temperature by a factor of two for the case of circular polarisation was observed for the ultra-thin targets. For the thicker targets, a decrease in the effective temperature of the hot electrons was also found for circular polarisation.

Figure 6.12: Optical transition radiation at $2\omega_0$, for linear (a) and circular (b) polarisation. The images were taken on the 8° incidence experiment for the $50 \mu\text{m}$ thick Au foils. (Figure courtesy of C. Bellei)

This difference between linear and circular polarisation is supported by optical transition radiation (OTR) measurements at $2\omega_0$ from the rear of the target, which show a decrease in intensity for circular polarisation, see figure 6.12 [171]. OTR is

the light that is emitted when a charged particle crosses the interface between two media with different dielectric properties. The change in the electromagnetic fields surrounding the charged particles, which is caused by the different dielectric constants of the two media, leads to the emission of radiation [172–174]. The spectrum of the radiation follows the Fourier transform of the envelope of the hot electron current. Therefore an optical signal at twice the laser frequency indicates that the electrons that are accelerated into the target are bunched at the same frequency. The signal for linear polarisation is a factor of 6 stronger than the one for circular polarisation.

One thing that is clear is that for circular polarisation, the high frequency component of the ponderomotive force (Lorentz force) is turned off, as expected, and therefore is likely to turn off $\mathbf{j} \times \mathbf{B}$ heating. This clearly results in a major change in the heating mechanism observed here.

6.4.2 Thickness scan

Independent of the laser polarisation the increase of the effective temperature with target thickness has been observed in three different experiments. One might think that this trend is due to an attenuation effect for thicker targets, however, if this were the case, one would expect a decrease in number for the lower energy electrons with increasing thickness. However, the increasing number (see figure 6.7) of electrons, especially at higher energies, suggests that it is not merely a range effect.

Comparing the electron spectra for the different experiments exposes some differences. On this account, the laser parameters for the solid target experiments are summarised in table 6.1, and a summary of the effective temperatures from the different cases is plotted in figure 6.13(a). For the experiment with 40° incidence (see section 6.3, figure 6.10(a)) the electron numbers and energies were the highest. The increased number of electrons and higher electron energies for this experiment, compared to the one at nearly normal incidence (see section 6.2, figure 6.6(a)), is probably due to the higher laser intensity. This can be seen more clearly in figure 6.13(b), which compares T_{eff}/U_p for the different experiments and suggests a simple

Parameter	8° linear	40°	45°
Focusing geometry	f/3	f/3	f/3
Focal spot diameter (μm)	9.5	8	5
Peak intensity (Wcm^{-2})	9×10^{19}	7.2×10^{20}	7.6×10^{20}
Peak a_0	8.2	24	25

Table 6.1: A summary of the laser parameters for the solid target experiments for linear polarisation.

scaling with a_0 .

For the experiment at 8° incidence, however, we have to assume a pre-pulse at a similar intensity level as the one at 40°. The pre-plasma would look similar in the two experiments, but the peak intensity is lower for the 8° experiment. Therefore the increased electron energies and number of electrons for the experiment at 40° are in agreement with intensity increase due to self-focusing and/or the pre-acceleration in the pre-plasma due to an underdense acceleration mechanism.

The increased number of electrons, compared to the 45° experiment (see section 6.3, figure 6.10(b)), could be due to the bigger spot size. The intensity on target was higher for the 45° experiment, but the spot size was smaller, and filamentation is more likely to occur. Filamentation leads to an effectively lower intensity in the single filaments, even though the total intensity in the laser might be higher. Another possible explanation is a change in pre-plasma.

Figure 6.13(b) shows a general trend of increasing temperature with increasing target thickness. In agreement with the discussion above, one can see that the temperatures for the 45° incidence, are not very high, and do not even approach the ponderomotive potential. The temperatures for the 40° incidence, at least for the shots for the 10, 20 and 100 μm thick targets seem to approach the ponderomotive potential, whereas for the shots at 25 and 50 μm , they are well above the ponderomotive potential. For the case of 8° incidence, the temperatures go to above the ponderomotive potential.

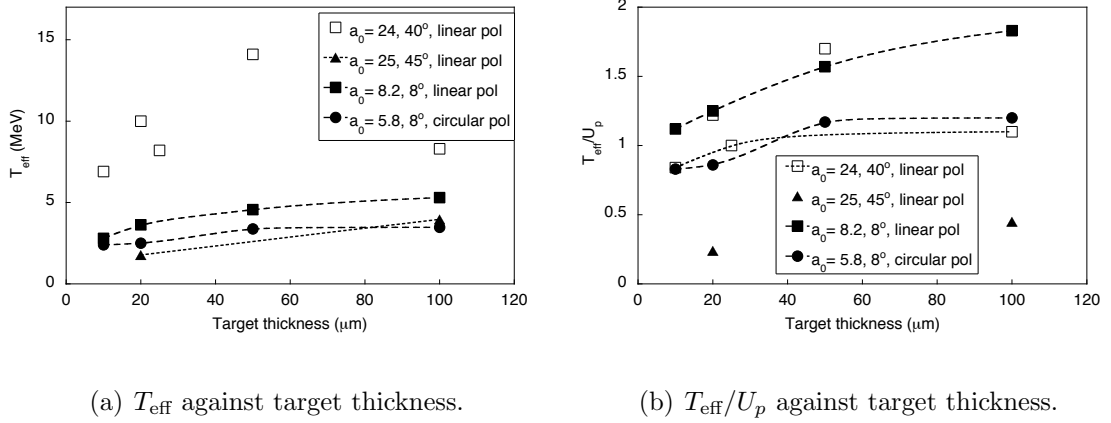


Figure 6.13: Plots of (a) T_{eff} and (b) T_{eff}/U_p against target thickness for different experiments. The angles indicate the angle of laser incidence onto the target. Note: The lines are only indications of the rough trend of the data.

6.5 Simulations

6.5.1 Polarisation comparison for thick targets

For the OSIRIS (see section 4.3) simulations presented here, a stationary box is used to observe the interaction of a laser pulse, $a_{0\text{lin}}=8.7$, with an overdense plasma. Gold has a solid density of 5.9×10^{22} atoms/cm³ ($\approx 59n_c$). And with intensities around 10^{20} Wcm⁻², for gold, a Au²⁰⁺ ionisation is expected. The plasma in the simulation has an electron density of $50n_c$ and is 10 μm thick (60grid cells). Here the electron density of $50n_c$ was chosen, to be able to resolve $c/\omega_p = 23.8$ nm within the computational constraint. The charge to mass ratio put into the simulation for the gold ions assumed full ionisation. The laser beam has a gaussian profile of 0.5 ps FWHM duration and the intensity at the focal spot has a FWHM of 8.5 μm . The computational grid has longitudinal (z) and transverse (x) dimensions of 84 μm by 28.5 μm , with resolutions of 8.4 nm ($\approx \lambda_p/18$) and 10.5 nm ($\approx \lambda_p/14$) in z - and x -direction respectively. The target is placed 3.5 μm into the box. The simulations were performed with linear and circular laser polarisations.

The electron density plots for three different times in the simulations are shown in figure 6.14. At 900 fs into the simulation (see figures 6.14(c) and 6.14(d)) the

deformation of the front surface, due to hole boring, is the deepest. After that time, the critical surface moves back towards the initial surface. The initial surface of the target is at $3.5 \mu\text{m}$, and the maximum amount that the critical surface moves back to is $5 \mu\text{m}$ for the case of circular polarisation. For the case of linear polarisation the critical surface is not that distinct. Therefore the hole boring is more efficient for circular polarisation, due to the constant ponderomotive push and reduced heating.

For circular polarisation one can see a higher number of electrons being accelerated at the surface of the “hole”. The dense electron front moves with a velocity of $(6.1 \pm 0.2) \times 10^6 \text{ ms}^{-1}$, which corresponds to an ion energy of around $\approx 37 \text{ MeV}$. This indicates shock acceleration. The observed velocity of the critical surface, or hole boring velocity, is $(2.2 \pm 0.5) \times 10^6 \text{ ms}^{-1}$. This is about a factor of three larger than the expected hole boring velocity* $0.8 \times 10^6 \text{ ms}^{-1}$. This difference could possibly be attributed to the lower density in the simulation.

Note that the ion density follows the electron density closely. Furthermore for circular polarisation, the sheath of electrons/ions that is accelerated is bigger than for linear polarisation. This is in agreement with the constant ponderomotive expulsion for circular polarisation, which leads to an easier accessibility of shock acceleration [36]. The pulse length however is too long and the target too thick to be able to access pure Radiation Pressure Acceleration (RPA), where the entire foil is accelerated uniformly [35].

* $v_s = \sqrt{\frac{2(1+R)I_0}{m_i n_i c}} = 0.8 \times 10^6 \text{ ms}^{-1}$ [175], where R is the reflectivity of the plasma, which is assumed to be 1, I_0 is the peak intensity of the laser, m_i is the ion mass, which for gold is $3.27 \times 10^{-25} \text{ kg}$, and n_i is the ion density, which for gold is $5.9 \times 10^{22} \text{ cm}^{-3}$.

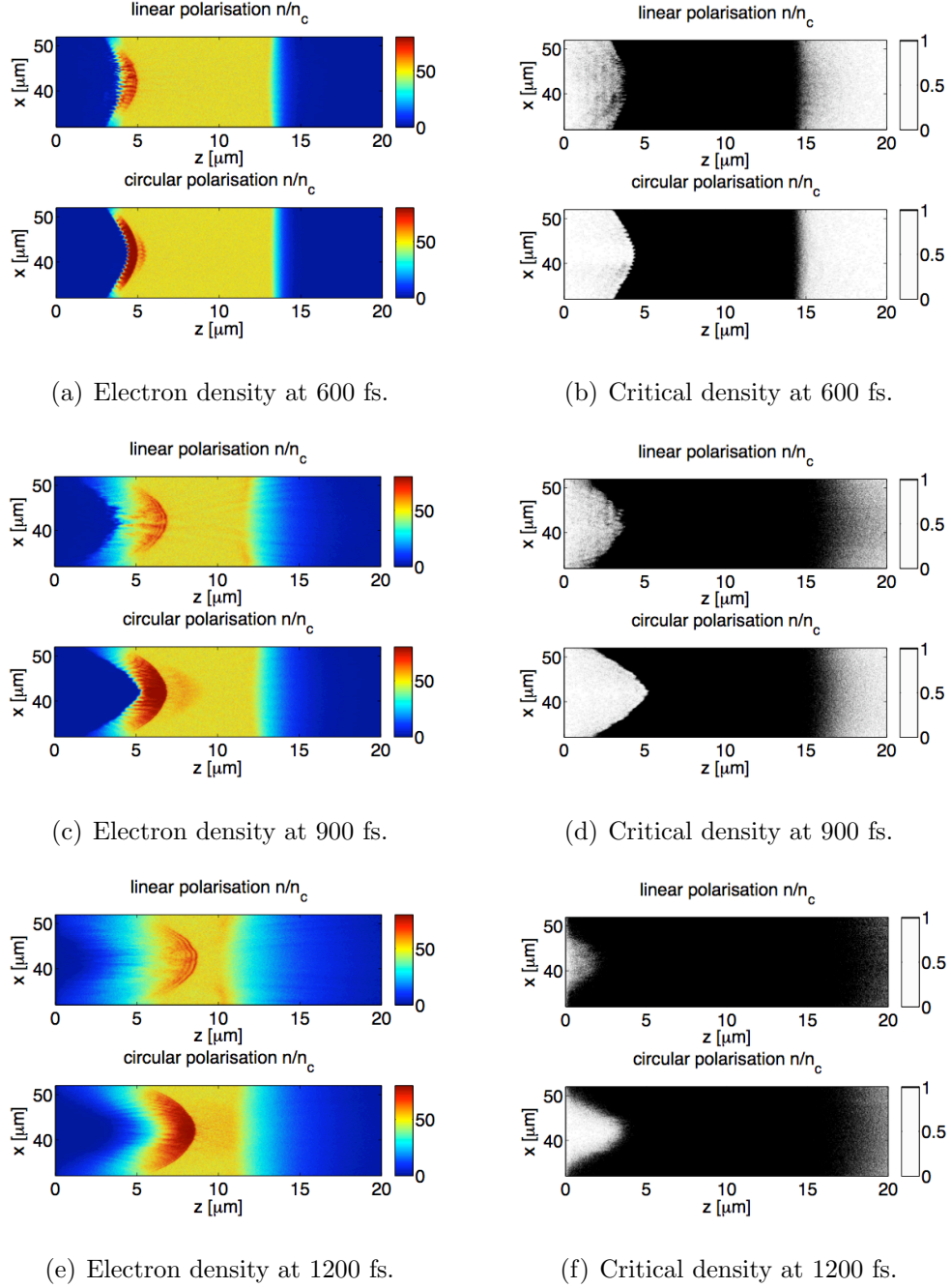


Figure 6.14: The electron density of a 10 μm thick target, for different time steps in the simulation. The plots on the left hand side show the electron density up to $80n_c$, and the right column shows the critical density boundaries.

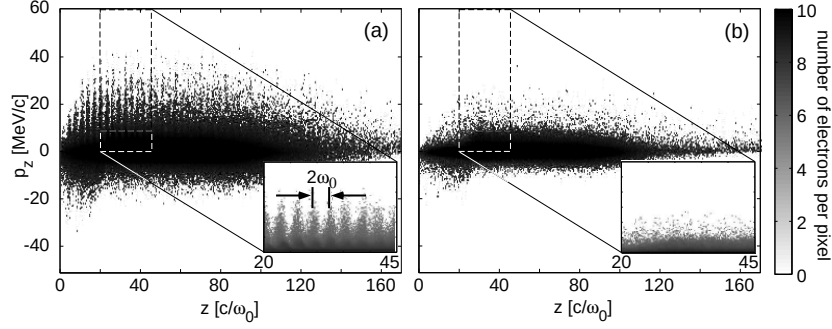


Figure 6.15: Phase space diagrams, for electrons in the laser direction at 738 fs into the interaction. p_z is the momentum in the laser direction (z). Linear polarisation (a) shows bunching at $2\omega_0$. This bunching cannot be observed for circular polarisation (b).

Figure 6.15 shows the phase space at 750 fs into the simulations, i.e. after the start of the laser pulse, for linear (a) and circular (b) polarisations. For the case of linear polarisation (with the electric field in the plane of the simulation) one can observe bunching of the electrons in the laser direction. The bunching of the electrons in the inset in figure 6.15(a) has a wavelength of 525 nm, which is exactly $\lambda_0/2$. This indicates $\mathbf{j} \times \mathbf{B}$ heating to be the dominant acceleration mechanism. For the case of circular polarisation (6.15(b)), the bunching of the electrons reduces, showing that the $\mathbf{j} \times \mathbf{B}$ heating is turned off. This reflects in the effective temperatures of the simulations in a similar way, as in the experimental data. At 750 fs into the simulation, we obtain a temperature of (1.5 ± 1) MeV for circular polarisation, and (3 ± 1) MeV for the linear polarised case. The reason why the bunching does not disappear completely for circular polarisation, is that due to hole boring, the interaction surface is not flat for the whole duration of the pulse, and the beam therefore develops an increasingly linear component as time progresses.

Figure 6.16 shows the electron spectra (in the forward direction) from the simulations, including all electrons in the box. The different colours correspond to different time steps in the simulations. Figure 6.16(a) shows the spectra for linear polarisation, and figure 6.16(b) the spectra for circular polarisation. For linear polarisation the electrons are heated faster than for circular polarisation (also see figure 6.17(a)). In general we observe a higher temperature for the linear polarisation, which is in

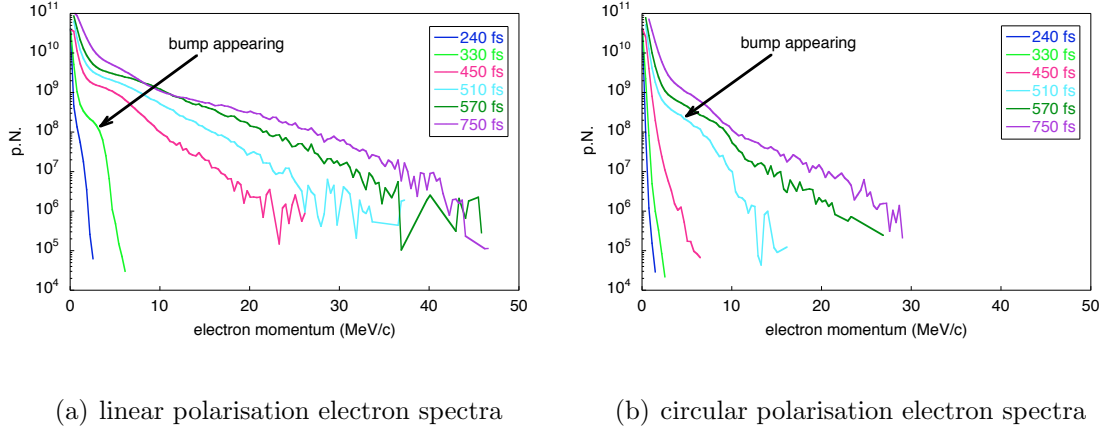


Figure 6.16: The electron spectra, for electrons moving in the laser direction, from the simulations of a 10 μm thick target and for different time steps in the simulation.

agreement with the experimentally observed data.

In the spectra in figure 6.16, one can make out a ‘bump’ at about ≤ 8 MeV for times later than 510 fs for circular polarisation and 330 fs for linear polarisation. This is more pronounced for linear polarisation, and washes out for later times. Looking at the phase space diagrams for this time, (330 fs, for linear polarisation) shows that the electrons in the target have gained up to 5 MeV, some higher. The electrons are bunched at $2\omega_0$, indicating $\mathbf{j} \times \mathbf{B}$ heating as the dominant acceleration mechanism. As the laser intensity increases with time, the energy of the bulk of accelerated electrons increases as well. For the case of circular polarisation the bump is not as pronounced, and it only becomes apparent for later times (between 450 and 510 fs). At the later times, even for circular polarisation, a slight modulation at $2\omega_0$ can be observed, indicating that the hole-boring has progressed so far that the circular polarised light, now has a linear component.

The time when hole boring disrupts RPA can, for circular polarisation, be estimated to be [176]:

$$\tau_{max} = \sqrt{\frac{2rcm_i n_i d}{I_0}},$$

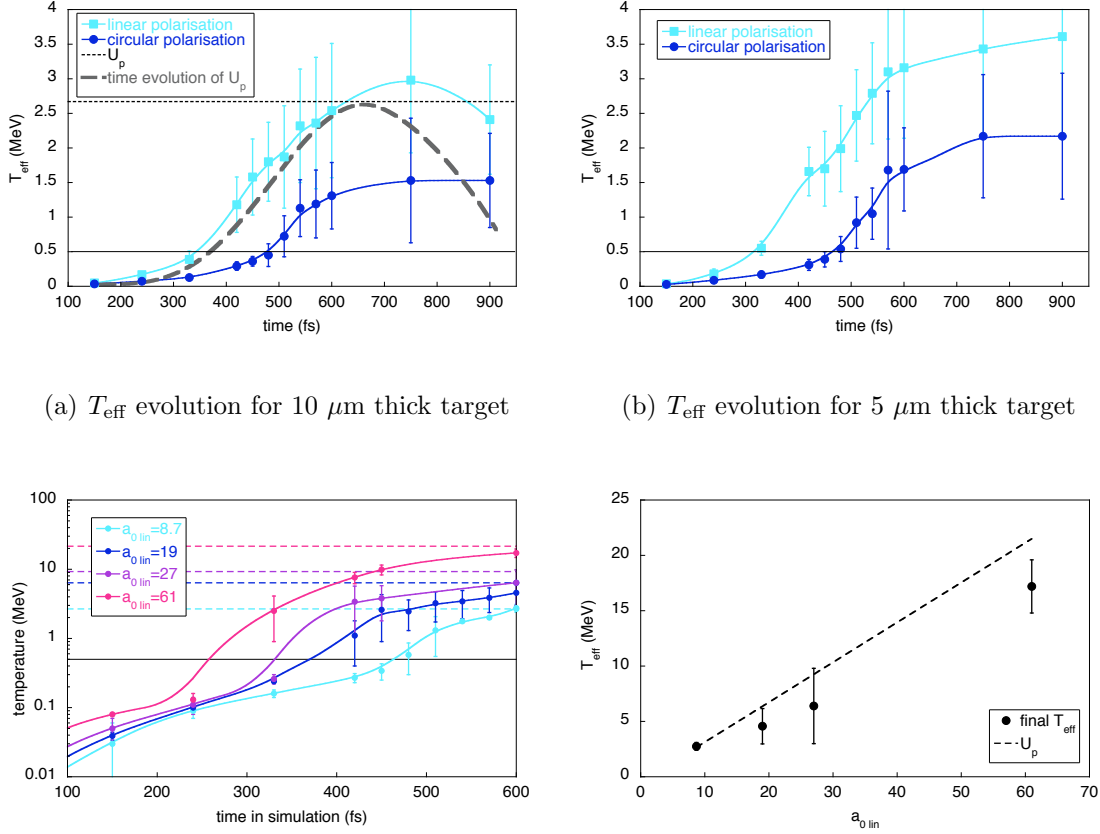
where r is the HWHM of the focal spot, m_i and n_i are the ion mass and ion density respectively, I_0 is the peak intensity of the laser pulse and $d = \sqrt{\frac{4\epsilon_0 I_0}{cn_e^2 e^2}}$. This

estimate assumes a target thickness d , which is a lot thinner than the one in the simulations and experiments. For the values from the simulation, $\tau_{max} = 162$ fs. For the parameters in the experiment, the estimated $\tau_{max} = 190$ fs.

Coming back to the temperature evolution for the two polarisations, which is plotted in figure 6.17(a), we can see that the time it takes to heat the electrons in the simulation to above 0.5 MeV is shorter for linear polarisation. For linear polarisation this time is 330 fs and for circular polarisation it is about 480 fs. This is in agreement with the fact that circular polarisation has constant ponderomotive expulsion which can sustain a large charge displacement [36]. Furthermore, the heating of the electrons is suppressed for circular polarisation. The time for the hole boring disruption is very close to the time difference in heating for linear and circular polarisation.

Simulations for a 5 μm thick target show a similar temperature evolution as for the 10 μm thick target, see figure 6.17(b). However, the temperatures increase faster and go up to higher values.

Further simulations, with the same setup, but using different intensities for circular polarisation, show that for increasing intensity, shorter laser pulses are desirable to prevent the heating of the targets, when trying to access the RPA regime. Figure 6.17(c) shows the temperature evolution for different intensities on target. It shows that for increasing laser intensity the heating of the target increases. The solid horizontal line in figure 6.17(c) corresponds to a temperature of 0.5 MeV, whereas the dashed lines indicate the ponderomotive potential, U_p , for the corresponding vector potentials. The intensities in the simulations range from $1 \times 10^{20} \text{ Wcm}^{-2}$ ($a_{0\text{lin}} = 8.7$) to $5 \times 10^{21} \text{ Wcm}^{-2}$ ($a_{0\text{lin}} = 61$). Figure 6.17(d) summarises the final temperatures for the different intensities, and compares them to the respective ponderomotive potential. One can see that the final temperature increases almost as U_p .



(a) T_{eff} evolution for 10 μm thick target (b) T_{eff} evolution for 5 μm thick target

(c) T_{eff} evolution for different intensities using circularly polarised light. The dashed lines indicate the respective U_p

(d) Final T_{eff} for different intensities using circularly polarised light. The dashed line indicates the corresponding U_p

Figure 6.17: Time evolution of T_{eff} , of the electrons moving in the laser direction, for (a) 10 μm and (b) 5 μm thick targets, comparing linear and circular polarisation. (c) compares the time evolution of the temperature in the simulations for different intensities using circularly polarised light for a 5 μm thick target. (d) shows the final temperature from (c) plotted against the corresponding a_0 , comparing it to U_p .

6.5.2 Thickness scan

As described in section 5.3, a stationary box was used for the OSIRIS (see section 4.3) simulations described here. The target thicknesses for these simulations are 10, 50 and 100 μm and the target is usually situated 3.5 μm ($21 (c/\omega_0)$) from the left boundary. The resolution is 8.4 nm in z and 10.5 nm in x . The laser propagates in the box from the left boundary to the right (z), and is polarised in the plane (linear polarisation). One time step corresponds to 15 fs, and the simulation is run with 4 particles per cell for both particle species (electrons and ions). A third test particle species is introduced, it has the same charge to mass ratio as the electrons in the simulation, but only 1 particle per cell, as well as different density profiles, to allow for tracking of this species, or better visualisation. The electron density is again set to correspond to $50n_c$ and the ion species is set up in the same way, as described in section 6.5.1. $a_{0\text{lin}}$ is set to 8.5, and the focal spot FWHM is 10 μm . The pulse shape is gaussian in time with a FWHM of 500 fs.

Figure 6.18 shows the density of the third species and the corresponding $p_z - z$ -plots for different time steps. In this simulation, the initial density profile of the “test” species is a thin slab at the target front surface of 0.168 μm thickness. The density of this test species was kept low ($\ll n_c$) to avoid it influencing the interaction. In figure 6.18 one can see the recirculation of the electrons at the rear surface of the target 6.18(d), this results in the increase of electron density at the target vacuum boundary and the electrons start to have a negative momentum shown in the $p_z - z$ diagram. At later times, the moving out of the rear surface can be observed 6.18(f).

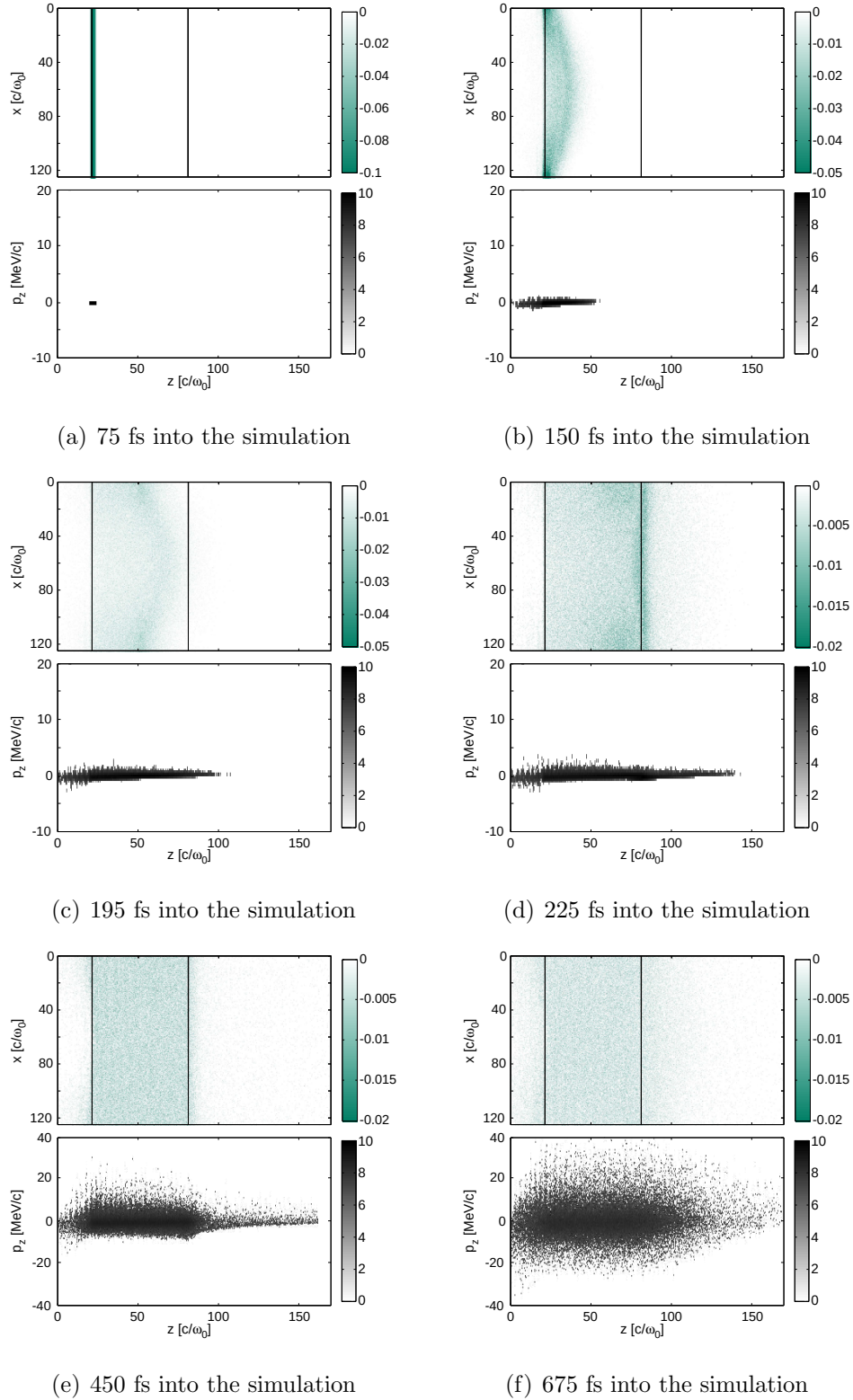


Figure 6.18: Density and phase space plots from a simulation with a 50 times critical density and $10 \mu\text{m}$ thick target. The species shown here is an electron species, that has an initial density profile as shown in 6.18(a) and has an initial density of $0.1n_c$. The black lines at 21 and 81 $z [c/\omega_0]$, indicate the target boundaries for the $50n_c$ electron species. The time steps were chosen to illustrate the initial recirculation of the electrons by the sheath field, shown in figure 6.20(a).

The two-dimensional longitudinal electric fields for the three thicknesses at two times after the sheath electric field starts growing are plotted in figure 6.19. The simulation box has an expansion in x , the transverse dimension, of $84 \mu\text{m}$. Assuming an electron beam divergence of 60° results in a theoretical divergence of $12 \mu\text{m}$ for the $10 \mu\text{m}$ thick target, $58 \mu\text{m}$ for the $50 \mu\text{m}$ thick target and $116 \mu\text{m}$ for the $100 \mu\text{m}$ thick target. One can see that the divergence of the electric sheath field for the $10 \mu\text{m}$ thick target is larger than $12 \mu\text{m} = 71 (c/\omega_0)$, and that the divergence for the $100 \mu\text{m}$ thick target is larger than the simulation box.

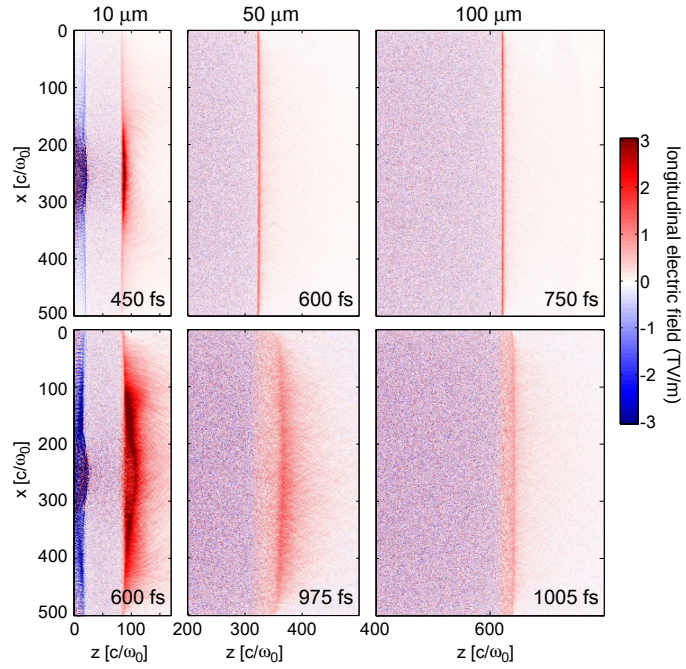


Figure 6.19: Longitudinal electric field for different target thicknesses and times in the simulations.

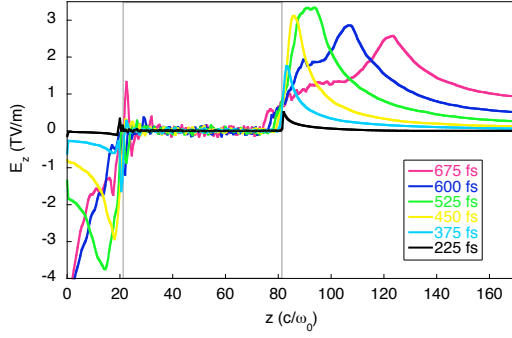
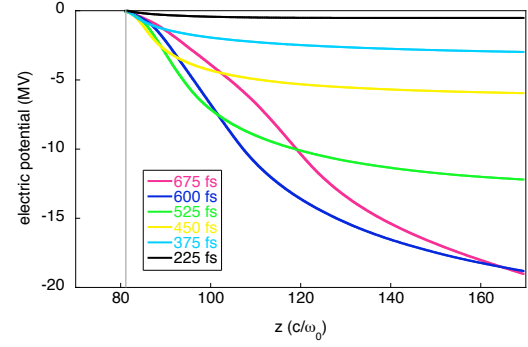
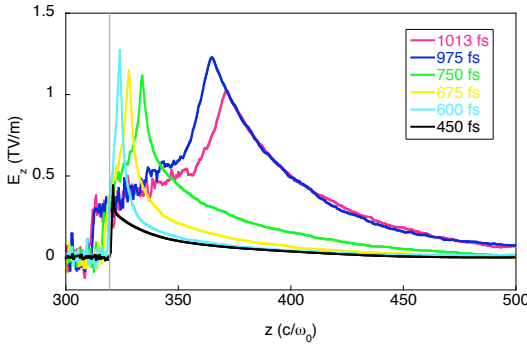
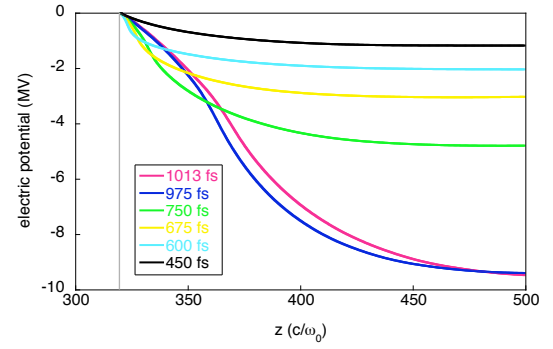
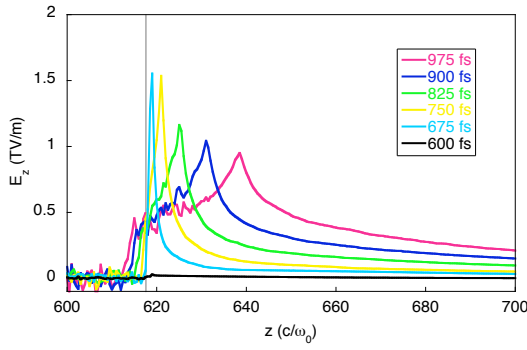
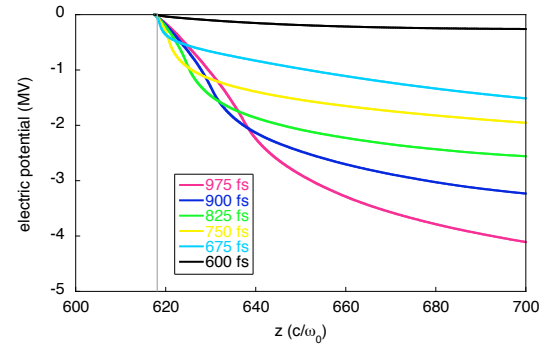
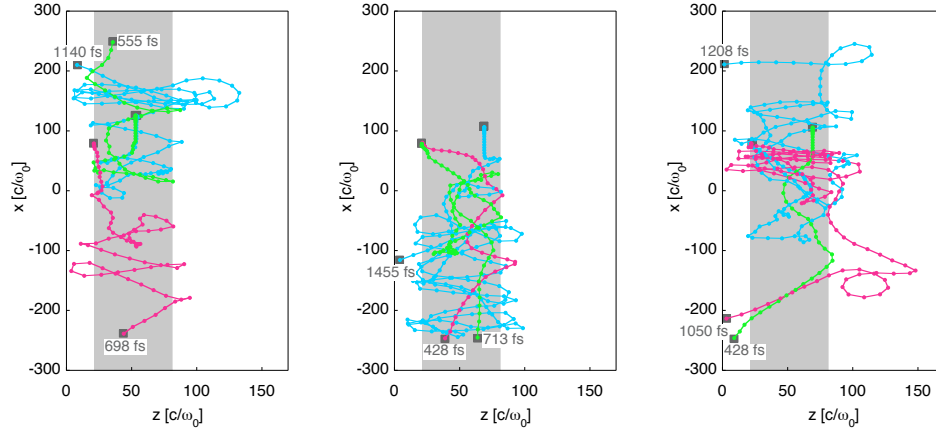
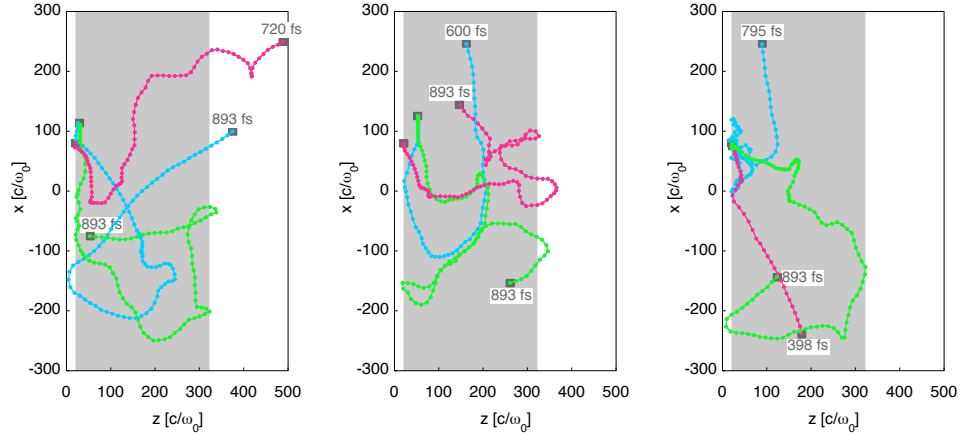
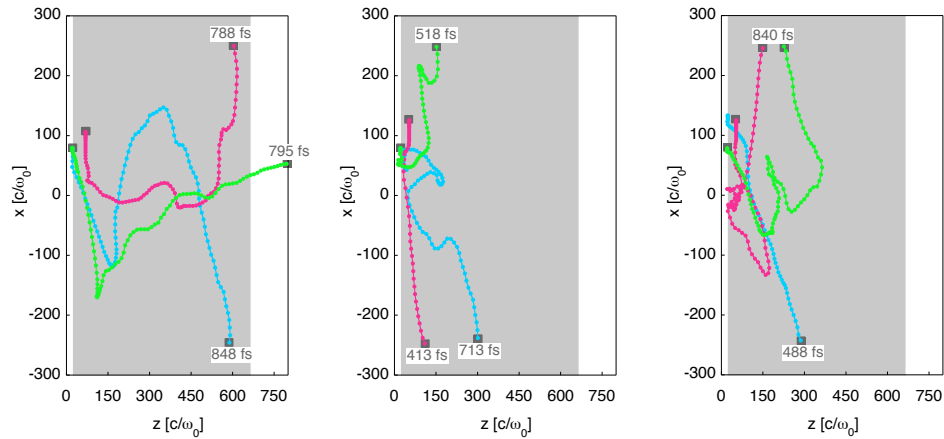

 (a) Electric field for 10 μm thick target

 (b) Electric potential for 10 μm thick target

 (c) Electric field for 50 μm thick target

 (d) Electric potential for 50 μm thick target

 (e) Electric field for 100 μm thick target

 (f) Electric potential for 100 μm thick target

Figure 6.20: The mean electric fields in laser propagation from the simulations for (a) 10 μm , (c) 50 μm and (e) 100 μm thick target and the electric potential for (b) 10 μm , (d) 50 μm and (f) 100 μm thick target. The grey lines indicate the initial target boundaries.

Figure 6.20 illustrates the evolution of the electric field in laser direction (z -direction), averaged in x over the simulation box size, for three different thicknesses, namely $10\ \mu\text{m}$ in figure 6.20(a), $50\ \mu\text{m}$ in figure 6.20(c) and $100\ \mu\text{m}$ in figure 6.20(e). Here the grey lines indicate the initial target boundaries. It can be seen that the field strength for the $10\ \mu\text{m}$ target is about 2.5 times higher than for the $50\ \mu\text{m}$ target. This is due to the fact that the electron beam has the same divergence for the two thicknesses, so that the area at the end of the target, from which electrons escape, is larger for thick targets than for the thin targets. If this were the only effect, one would expect a difference in field strength of a factor 25. However, for thin targets the area can be increased due to recirculation of the electrons, which reduces the difference in the strength of the sheath field. Due to the higher potential electrons trying to escape the thin target lose more energy trying to overcome the electric field.

The electric potential resulting from these electric fields is plotted in figures 6.20(b), 6.20(d) and 6.20(f) for the $10\ \mu\text{m}$, $50\ \mu\text{m}$ and $100\ \mu\text{m}$ thick targets, respectively. This shows that electrons trying to escape the $10\ \mu\text{m}$ thick target will lose more energy than the ones escaping the $50\ \mu\text{m}$ thick targets. For the two different target thicknesses, the difference in maximum electric potential is $\approx 10\ \text{MV}$.

Figure 6.21 shows different electron trajectories from the start until they leave the box (times indicated in plots) for a $10\ \mu\text{m}$ thick target, and linear laser polarisation. In order to be able to track the particles the density profile of the test species had a gradient in both directions. The particles were picked at random, and most of them recirculate due to the sheath electric field. The same plots, but for a $50\ \mu\text{m}$ and a $100\ \mu\text{m}$ target can be seen in figures 6.22 and 6.23, respectively. For the $10\ \mu\text{m}$ target, figure 6.21, numerous circulations can be seen, whereas for the very thick $100\ \mu\text{m}$ target, figure 6.21, this cannot be observed for the time the laser is active.


 Figure 6.21: Particle trajectories from simulation with 10 μm thick target.

 Figure 6.22: Particle trajectories from simulation with 50 μm thick target.

 Figure 6.23: Particle trajectories from simulation with 100 μm thick target. The grey areas indicate the initial target position. The time between steps is 7.5 fs. Note that here, the zero position on the y -axis is the centre of the simulation, whereas for the other simulation plots, that is not the case.

The time evolution of the forward momentum and the position in z , for six different particles for each target thickness, are plotted in figures 6.24 to 6.26. Again the recirculation of the electrons can be seen for the thin targets, as well as the increase of forward momentum at the front surface. Usually, for all of the different thicknesses, the electrons gain the most energy at the front surface closest to the centre of the focus.

Furthermore, for the linear polarisation electrons are accelerated to above the ponderomotive potential of $U_p = 2.6$ MeV. For the $10\text{ }\mu\text{m}$ thick target, the momentum an electron gains in the different recirculations appears to depend on the laser intensity at that time. The momentum gain for thicker targets, at least for the electrons plotted here is still above U_p , but is smaller compared to the thin target. This could be due to the tracked particles all originating from the front of the target, where the electrons gain momentum from the low intensity part of the laser. For the $100\text{ }\mu\text{m}$ thick target, the green trajectory in the right picture in figure 6.26 looks as if it is just turning around in the middle of the target. The corresponding trajectory is the green line in the third picture in figure 6.23, which shows that the turning around is more due to a transverse motion of the electron. Since the PIC code has no collisions, this is due to anomalous scattering.

Figure 6.27 and 6.28 show the electron trajectories and the time evolution of the forward momentum and the position in z for a $10\text{ }\mu\text{m}$ thick target with circular laser polarisation. Here the electrons are accelerated to lower energies than in the linear polarisation case. For the case of linear polarisation, the electrons gain momentum above U_p , whereas for the circular polarisation case, the energy electron gain in the laser direction corresponds to $U_p = 2.6$ MeV.

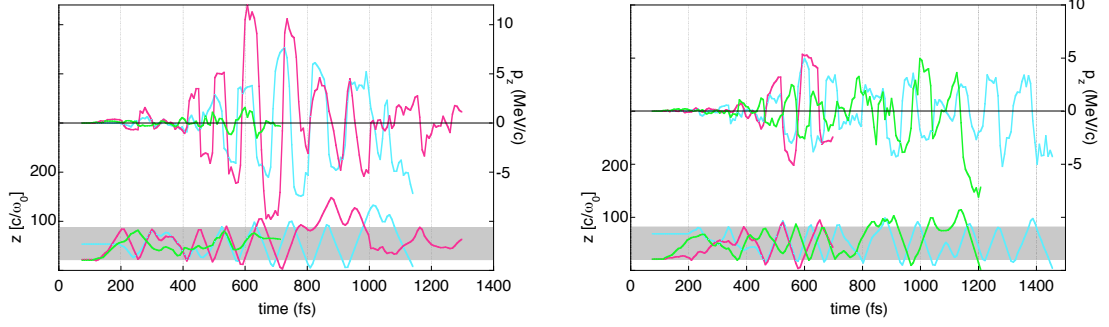


Figure 6.24: Particle position in z direction and momentum plotted against time in the simulation, for $10 \mu\text{m}$ thick target.

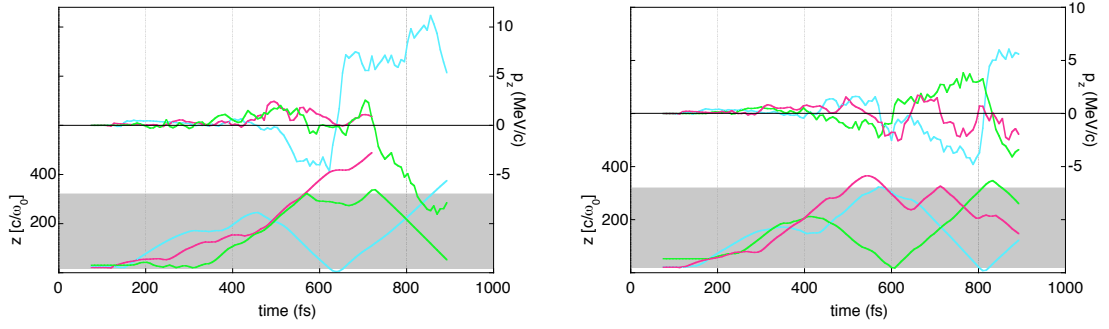


Figure 6.25: Particle position in z direction and the momentum plotted against time in the simulation, for $50 \mu\text{m}$ thick target.

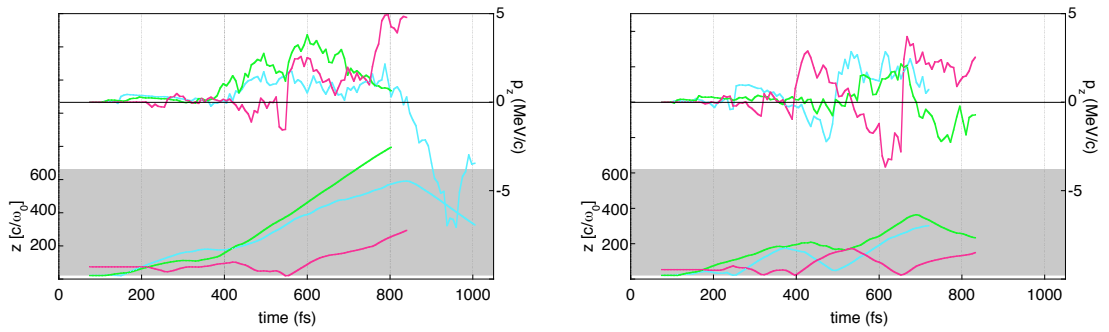


Figure 6.26: Particle position in z direction (left axis and bottom), and the momentum of the particle in z (right axis and top), plotted against time in the simulation, for $100 \mu\text{m}$ thick target.

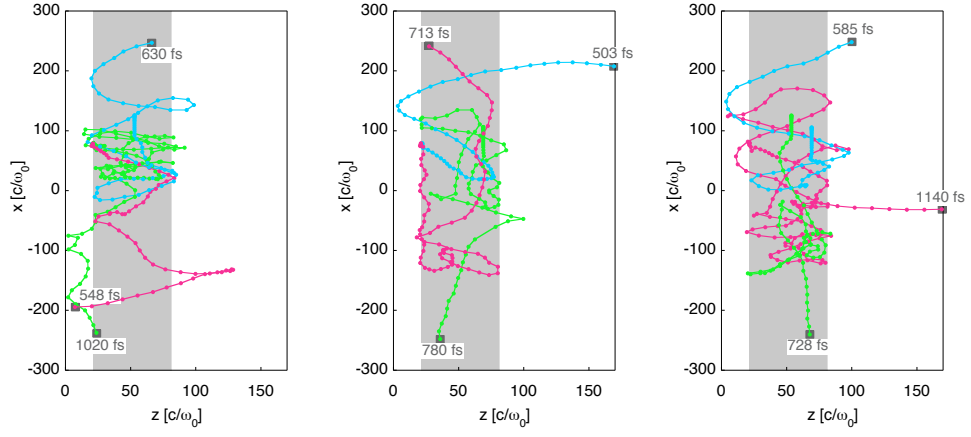


Figure 6.27: Particle trajectories from simulation with $10\ \mu\text{m}$ thick target for circular polarisation.

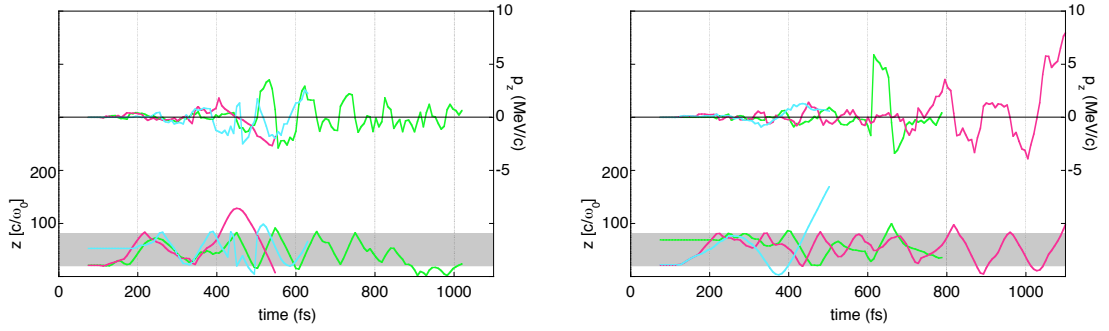


Figure 6.28: Particle position in z direction and momentum plotted against time in the simulation, for $10\ \mu\text{m}$ thick target and circular polarisation.

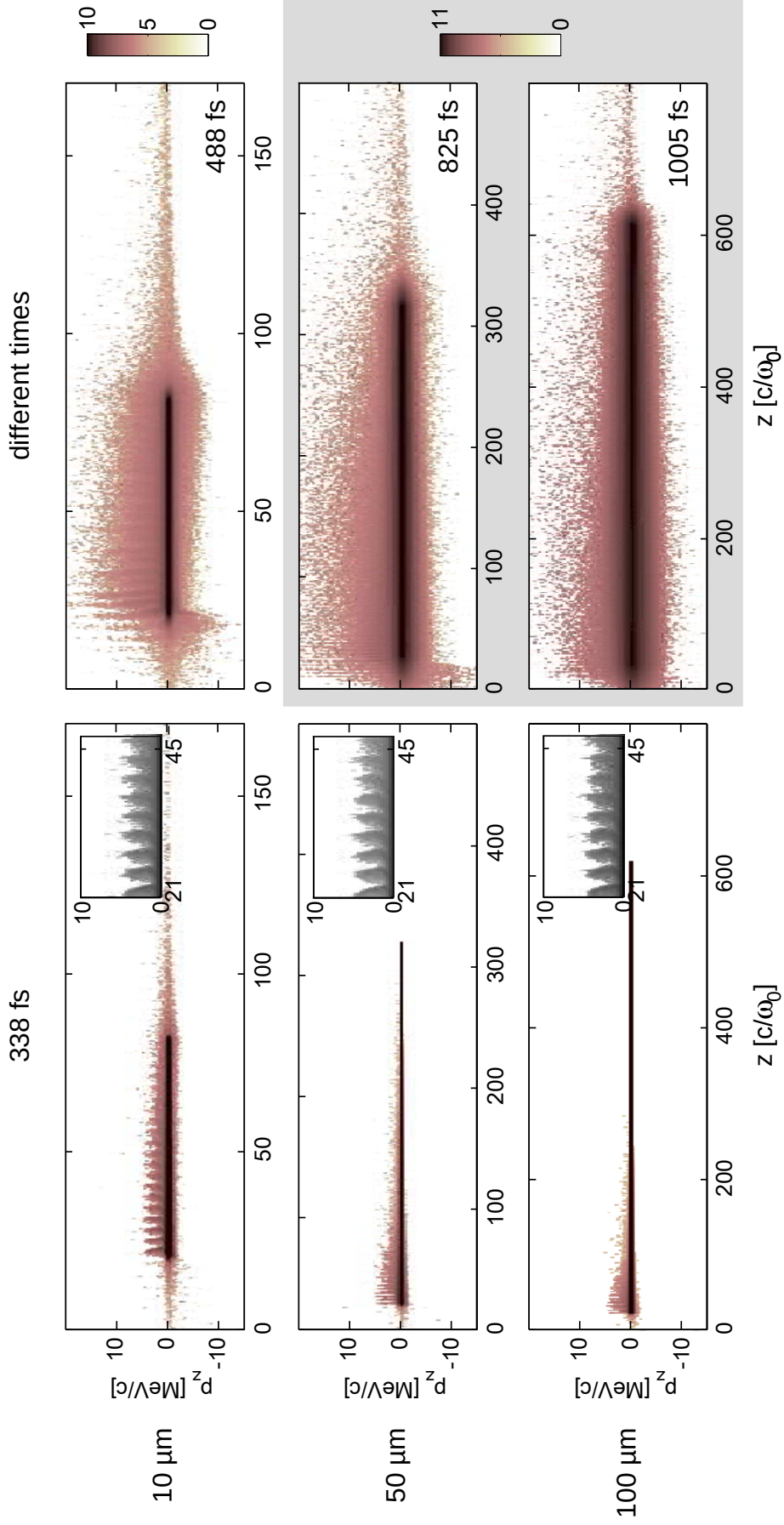


Figure 6.29: Comparison for the time evolution of the phase space in the simulations for different target thicknesses (rows) at two different times (columns). All targets start at $21 c/\omega_0$, and end at, $81 c/\omega_0$ ($10 \mu\text{m}$), $320 c/\omega_0$ ($50 \mu\text{m}$) and $618 c/\omega_0$ ($100 \mu\text{m}$).

In figure 6.29 the phase space for different target thicknesses (10, 50 and 100 μm , rows) is plotted for two different times in the simulations. The heating of the targets can be observed. When the 10 μm thick target is heated through, at about 338 fs, the thicker targets are heated over only half their thickness, or even less. One can see the bunching due to $\mathbf{j} \times \mathbf{B}$ heating in all targets. It is harder to see in this plot for the thicker targets, because of the increased plotting range in z . Therefore the insets show a zoom into the relevant areas. There are also returning electrons at all thicknesses, due to recirculation and in the case of thicker targets, due to anomalous scattering.

The right hand side of figure 6.29 shows the phase space at different times for the different simulations. These times correspond to roughly the same times after electrons have exited the target. For all target thicknesses, electrons are also traveling backwards, due to being stopped by the sheath field.

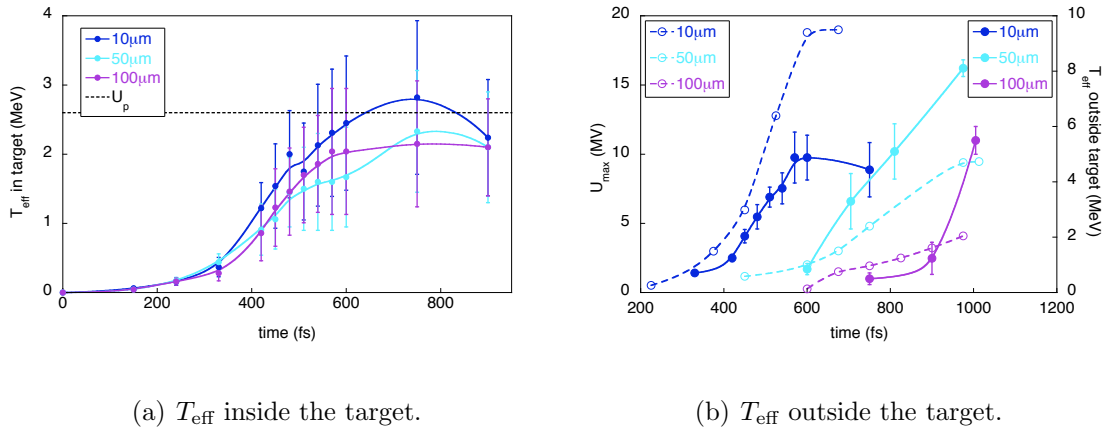


Figure 6.30: Time evolution of T_{eff} for different thicknesses (a) inside and (b) outside the target (solid lines). (b) also shows the time evolution of the maximum electric potential, due to the sheath field (dashed lines).

The time evolution of the effective temperatures inside the target for different targets, as well as the ponderomotive potential are plotted in figure 6.30(a). Typically the temperatures within the target approach a value slightly below the ponderomotive potential. For the 10 μm thick target, at some point, the temperature even exceeds the ponderomotive potential, this is probably due to recirculation.

The temperatures for the electrons outside the target are harder to obtain. Additionally for the thickest target, the simulation did not run long enough, due to temporal limitations on the computer cluster, for the electrons to see the maximum potential. However, for the times at which the maximum potential is reached, for the 10 μm and the 50 μm thick target, we obtain electron temperatures of (5.5 ± 0.1) MeV and (8.0 ± 0.3) MeV, respectively. This can be seen in figure 6.30(b), where the time evolution of the maximum electric potential due to the sheath field is plotted, as well as the temperature of the electrons outside the target. The decrease in temperature for the 10 μm thick target at late times (750 fs) is most likely due to energetic electrons leaving the stationary simulation box. Compared to the experiment, the temperatures from the simulation are higher, but show the same increase with increasing target thickness.

Since the temperatures within the targets approach the same value it is apparent that independent of target thickness the electric potential at the boundary causes the observed thickness dependence of the experimentally measured temperatures (see section 6.4.2).

Chapter 7

Summary

The main focus of this thesis was the investigation of electron acceleration from ultra-high intensity laser interactions with plasmas. These hot electrons are responsible for numerous other mechanisms, such as x-ray production and ion acceleration, and a detailed understanding of the electron acceleration is therefore important. The results from the experimental chapters are summarised below. They are split into electron acceleration from underdense and overdense plasmas.

7.1 Electron acceleration from underdense plasmas

In chapter 4, intensity scans, density scans, f-number comparisons as well as interaction length scans were performed experimentally and simulations were conducted to try and help understand the experimental results better. The chapter was split into three sections. Two experimental sections investigated the intensity regimes above and below $1 \times 10^{20} \text{ Wcm}^{-2}$, and the third section discussed PIC simulations.

Intensity scans above $1 \times 10^{20} \text{ Wcm}^{-2}$ show the scaling of the electron temperature with a_0 , $T_{\text{eff}} \propto \sqrt{I} \propto a_0$. However, for low a_0 the scaling breaks down, indicating that a different mechanism is responsible for the acceleration. As the intensity decreases, the shape of the electron spectra also becomes less maxwellian. In the relevant simulations three areas of acceleration are observed at the same time.

Their relative importance depends on a_0 . In the case of high a_0 electrons are accelerated to above the dephasing energy, and therefore a high maximum electron energy as well as temperature, indicating that DLA is dominant. Whereas for the case of low a_0 , the maximum energy electrons are accelerated to is the dephasing energy, which suggests that the SMLWFA regime takes over. This is also supported by x-ray measurements showing a dramatic decrease of radiation for the lowest intensity.

The experimental results also show a density dependence for the efficiency of the electron acceleration, see figures 4.7, 4.8 and 4.9. And for a modest change in f-number ($f/5$ is only 1.7 times larger than $f/3$) we observe a significant improvement of the electron acceleration, in electron number and maximum energy, in the case of $f/5$. For both f-numbers an inverse density dependence of T_{eff} , maximum energy and number of electrons, can be observed for high densities. However, the lowest density does not follow this scaling. $f/5$ -focusing shows better acceleration than $f/3$ -focusing, even though it features a lower intensity. But it is also expected to lead to a longer Rayleigh length, z_R , indicating that the interaction length is longer for $f/5$ -focusing. In both simulation and experiment, the optimum density increases for smaller f-numbers.

Additional beam divergence measurements show a constant divergence for $f/5$ -focusing, whereas $f/3$ -focusing shows decreasing divergence. This suggests that in the case of $f/5$ -focusing the electrons pick up radial as well as longitudinal momentum, supporting the assumption that higher energies are accelerated by DLA.

Interaction length scans were conducted at two different intensities, namely $1 \times 10^{20} \text{ Wcm}^{-2}$ and $1.6 \times 10^{21} \text{ Wcm}^{-2}$. For the lower intensity the effect of thermalisation with increasing interaction length was observed, see figure 4.14. At the higher intensity, the interaction length does not play as important a role as for lower intensities and the main acceleration of the electrons seems to happen within the 1st mm. The electron spectra in this experiment feature a two-temperature distribution, which could be due to two DLA mechanisms, since even the low temperature part extends to above the dephasing energy.

These measurements are relevant for future laser plasma based x-ray sources and particle accelerators. As an extension of this work, experiments fostering a better

understanding of the transition between the two regimes (DLA and SMLWFA), and extensions to higher a_0 would be interesting. In particular as we have shown, that the bubble size is a_0 dependent. These studies are likely to happen with the 10 PW upgrade to VULCAN. Ideally, the f-number comparison could be expanded to matching of the focal spot size to the plasma wavelength for this “long” pulse regime, to see if the efficiency increases further for larger f-numbers.

Comparing different f-numbers at intensities below $1 \times 10^{20} \text{ Wcm}^{-2}$, shows that larger f-numbers lead to higher electron numbers and maximum electron energies, even if the peak intensity for these large f-number shots is more than a factor of ten lower than for the shorter focal lengths, see table 4.2. Here the interaction length plays an important role in the thermalisation of the electrons, and therefore changes the spectra significantly.

7.2 Electron acceleration from overdense plasmas

Chapters 5 and 6 investigated several issues in the overdense regime. One of them was the effect of different laser contrast ratios on the escaping electrons. This showed a higher number of electrons and electron temperature for the low contrast ratio (high pre-pulse level), perhaps indicating that underdense mechanisms influence the hot electron generation. As an extension to the experimental results, the pre-plasma characterisation and simulations with the measured scale lengths are under way.

We have also studied the electron acceleration from gold foils ranging from thicknesses of 10 to 100 μm , using different laser polarisations. Two major trends in the effective temperatures of the escaped electrons are observed (see figure 6.8), namely a polarisation and a thickness dependence.

The polarisation dependence, which has been seen in two independent experiments, described in chapters 5 and 6.2, can be attributed to the turn off, or reduction, of the $\mathbf{j} \times \mathbf{B}$ heating for the case of circular polarisation, independent of target thickness. In the experiment with close to target normal incidence, the circular polarisation is not purely circular, but rather elliptical, since we are hitting the

target at a slight angle. This effect is further enhanced by hole boring, and could explain the small difference between linear and “circular” polarisation. The OTR measurements, shown in section 6.4.1, also support this, as the emission at $2\omega_0$ does not disappear completely for circular polarisation. Simulations additionally show a linear dependence of the effective temperature on a_0 for circular polarisation, that follows the scaling of the ponderomotive potential closely.

The second trend is the thickness dependence of the effective temperature – which has been observed in three experiments, described in sections 6.2 and 6.3. We observe a decrease of the temperature with decreasing target thickness, independent of laser polarisations. Simulations show an effective temperature inside the target that is independent of target thickness, which indicates that the thickness dependence of the experimentally measured electrons is due to the energy loss in the rear surface sheath field. This result, that the temperature dependence of the escaping electrons is due to the electric sheath field, implies that the temperature inside the target cannot be inferred simply from the escaped electrons.

In general, controlling the electron energy and temperature is important for ICF, in particular the fast ignition scheme, as it allows for the possibility of selective energy deposition at the core of the pre-compressed fuel pellet. The investigation of hole boring for circular polarisation is also important for the ICF scheme. In particular circular polarisation and low pre-pulses may help limit the temperature of the fast electrons generated in either fast or shock ignition schemes. Therefore increasing the effectiveness of both schemes.

The experimental results on the polarisation dependence are also promising, in the view of the fact that, for the new generation of short pulse PW class laser systems, it might be possible to access the RPA regime using circular polarisation and ultra-thin targets.

Frequently used symbols

A	vector potential	λ_μ	laser wavelength in μm
a	normalised vector potential	m	mass/electron mass
a_0	amplitude of a	m_e	electron mass
B	magnetic field	m_i	ion mass
c	speed of light	n	density
$c\tau$	laser pulse length	n_0	ambient electron density
d	FWHM of focal spot	n_c	critical density
e	electron charge	n_e	electron density
$\mathbf{e}_{x,y,z}$	unit vector	n_i	ion density
E	electric field	η_R	refractive index
E	electric field strength	w_0, σ_0	waist of focal spot
$\mathcal{E}_L$	laser energy	ω	laser frequency
$\mathcal{E}_{Laser}$	laser energy	ω_0	laser frequency
$\mathcal{E}_i$	ionisation energy	ω_β	betatron frequency
ϵ_0	permittivity of free space	ω_p	plasma frequency
F	focusing force due to electric and magnetic fields	P	laser pulse power
F_p	ponderomotive force	P_c	critical power
$f_\#$	f-number	p	momentum
f_{number}	f-number	ϕ	scalar potential
Γ	Keldysh parameter	q	charge
γ	Lorentz factor	r_L	Larmor radius
γ_e	betatron resonance energy	ρ	charge density
I	laser intensity	T_{eff}	effective temperature
I_0	laser intensity	τ	laser duration
I_{18}	laser intensity in 10^{18}Wcm^{-2}	U_p	ponderomotive potential
k	wave number	v	velocity
k_B	Boltzmann constant	v_g	group velocity
L	Lagrangian/density scale length	v_{ph}	phase velocity
L	target thickness	v_{os}	a_0c
L_{pd}	laser depletion length	z_R	Raleigh length
L_d	dephasing length	$\langle \rangle$	time average
λ	laser wavelength	$\perp$	subscript perpendicular to laser propagation
λ_0	laser wavelength	$\parallel$	subscript parallel to laser propagation
λ_p	plasma wavelength		

Fundamental constants and quantity conversions

Symbol	Name	Quantity	Unit
e	elementary charge	1.602×10^{-19}	C
u	unified atomic mass unit	1.66×10^{-27}	kg
m_e	rest mass of an electron	9.11×10^{-31}	kg
m_p	rest mass of a proton	1.67×10^{-27}	kg
ϵ_0	permittivity of free space (electric constant)	8.85×10^{-12}	Fm ⁻¹
μ_0	permeability of free space (magnetic constant)	$4\pi \times 10^{-7}$	
c	speed of light in vacuum	2.998×10^8	ms ⁻¹
h	Planck's constant	6.63×10^{-34}	Js
k_B	Boltzmann's constant	1.38×10^{-23}	JK ⁻¹
$m_e c^2$	energy equivalence of electron rest mass	0.511	MeV
a_0	peak vector potential	$0.86\lambda_\mu\sqrt{I_{18}}$	
ω_p	plasma frequency	$56400\sqrt{n_{e[\text{cm}^{-3}]}}$	s ⁻¹

Appendix A

A.1 Interaction with single atom and electron

Starting with Maxwell's equations:

$$\nabla \times \mathbf{B} = \mu_0 \mathbf{j} + \epsilon_0 \mu_0 \frac{\partial}{\partial t} \mathbf{E} \quad (\text{Ampère's circuital law, with Maxwell's corrections}) \quad (\text{A.1})$$

$$\nabla \times \mathbf{E} = -\frac{\partial}{\partial t} \mathbf{B} \quad (\text{Faraday's law of induction}) \quad (\text{A.2})$$

$$\nabla \cdot \mathbf{E} = \rho / \epsilon_0 \quad (\text{Gauss's law}) \quad (\text{A.3})$$

$$\nabla \cdot \mathbf{B} = 0 \quad (\text{Gauss's law for magnetism}) \quad (\text{A.4})$$

With the vector potential $\mathbf{A}$, and the scalar potential ϕ , the electric and magnetic fields can be expressed as:

$$\mathbf{E} = -\frac{\partial \mathbf{A}}{\partial t} - \nabla \phi \quad (\text{A.5})$$

$$\mathbf{B} = \nabla \times \mathbf{A} \quad (\text{A.6})$$

The relativistic Lagrangian of a particle with charge q moving in such a potential is given by:

$$L(\mathbf{r}, \mathbf{v}, t) = -mc^2 \sqrt{1 - \frac{v^2}{c^2}} + q\mathbf{v} \cdot \mathbf{A} - q\phi \quad (\text{A.7})$$

Using the Euler–Lagrange equation,

$$\frac{d}{dt} \frac{\partial L}{\partial \mathbf{v}} - \frac{\partial L}{\partial \mathbf{r}} = 0 \quad (\text{A.8})$$

$$\frac{d}{dt} (\mathbf{p} + q\mathbf{A}) = q(\nabla \mathbf{A}) \cdot \mathbf{v} - \nabla q\phi \quad (\text{A.9})$$

By using the fact that $\frac{d}{dt} = \frac{\partial}{\partial t} + \mathbf{v} \cdot \nabla$ and the vector identity $\mathbf{v} \times \nabla \times \mathbf{A} = (\nabla \mathbf{A}) \cdot \mathbf{v} - (\mathbf{v} \cdot \nabla) \cdot \mathbf{A}$ we can transform the equation of motion (A.9) into:

$$\frac{d\mathbf{p}}{dt} = q(\mathbf{E} + \mathbf{v} \times \mathbf{B}) \quad (\text{A.10})$$

Assuming a free electron in a plane electromagnetic wave propagating in the z -direction and polarised in the x -direction. In the nonlinear relativistic regime, the magnetic field component becomes non-negligible and curves the electron trajectory towards the z -direction, (compared to the linear regime, where the electron oscillates in the perpendicular (xy) plane under the action of the electric field). As a result the electron has a longitudinal as well as transverse motion with respect to the light propagation direction. Further assuming that the scalar potential $\phi = 0$, and the vector potential is a travelling function such that $\mathbf{A} = A(z - ct)\hat{\mathbf{x}}$.

Since $\mathbf{A}$ does not vary in x and y :

$$\Rightarrow \frac{d}{dt} \frac{\partial L}{\partial v_x} = 0 \quad (\text{A.11})$$

Since $\partial L / \partial \mathbf{v} = m\gamma \mathbf{v} + q\mathbf{A}$ is the canonical momentum, equation (A.11) states that this is conserved, i.e. $p_x(t) + qA_x(t) = p_x(0) + qA_x(0)$. For an electron at rest before the arrival of the laser, this reduces to $p_x = -qA_x$. From the Hamiltonian principal the energy is given by:

$$\frac{d\mathcal{E}}{dt} = -\frac{\partial L}{\partial t} \quad (\text{A.12})$$

Because $\mathbf{A} = f(z - ct)$, this can be rewritten as:

$$\frac{d\mathcal{E}}{dt} = c \frac{\partial L}{\partial z} \quad (\text{A.13})$$

Using the z -component of the Euler-Lagrange equation (A.8):

$$\frac{d\mathcal{E}}{dt} = c \frac{\partial L}{\partial z} = c \frac{d}{dt} \frac{\partial L}{\partial v_z} = c \frac{dp_z}{dt} \quad (\text{A.14})$$

This gives the second constant of motion:

$$\frac{d}{dt}(\mathcal{E} - cp_z) = 0 \quad (\text{A.15})$$

or $\mathcal{E}(t) - cp_z(t) = \mathcal{E}(0) - cp_z(0)$. Since the total energy of a particle is $\mathcal{E} = \gamma mc^2$ and $\gamma = \sqrt{1 + p_x^2/(m_e c)^2 + p_z^2/(m_e c)^2}$ then it follows for a particle with zero initial momentum*:

$$p_z = \frac{p_x^2}{2mc} = \frac{(qA_x)^2}{2mc} = mc \frac{a^2}{2} \quad (\text{A.16})$$

*For zero initial momentum: $\mathcal{E}(t) - \mathcal{E}(0) = cp_z(t)$, and $(\gamma - 1)m_e c^2 = cp_z(t)$. Therefore $\gamma = \frac{p_z(t)}{m_e c} + 1$. Setting this equal to $\gamma = \sqrt{1 + p_x^2/(m_e c)^2 + p_z^2/(m_e c)^2}$ gives equation (A.16).

where the normalised vector potential $a = qA/mc$ was defined. Equation (A.16) describes a parabolic trajectory in momentum space.

To derive the trajectory of the electron in space we need to integrate the equations of motion. In the following we use a normalised set of variables where momenta are normalised to mc and velocities are normalised to c . For completeness let us assume that $\mathbf{a} = (a_x, a_y, 0)$. The constants of the motion then produce the following:

$$\begin{aligned} p_x &= a_x \\ p_y &= a_y \\ \gamma &= p_z + 1 = 1 + \frac{a^2}{2} \end{aligned} \tag{A.17}$$

Since $p = \gamma \mathbf{v}/c$ in the normalised units we can write:

$$\begin{aligned} \gamma \frac{dx}{dt} &= ca_x \\ \gamma \frac{dy}{dt} &= ca_y \\ \gamma \frac{dz}{dt} &= c \frac{a^2}{2} \end{aligned} \tag{A.18}$$

Now we substitute variable $\tau = t - z(t)/c$, i.e. to a variable which describes the phase of the wave since $a = a(t - z/c)$. All of the equations include the operator $\gamma \frac{d}{dt}$. We need to turn this into a differential in terms of τ .

$$\gamma \frac{d}{dt} = \gamma \frac{d\tau}{dt} \frac{d}{d\tau} = \gamma \left(1 - \frac{dz}{c dt} \right) \frac{d}{d\tau} = \left(1 + \frac{a^2}{2} - \frac{c a^2}{c 2} \right) \frac{d}{d\tau} = \frac{d}{d\tau} \tag{A.19}$$

This reveals that not only does τ relate to the phase of the wave but it is also the proper time of the particle $\tau = t/\gamma$. We can now rewrite equations (A.18) as:

$$\begin{aligned} \frac{dx}{d\tau} &= ca_x \\ \frac{dy}{d\tau} &= ca_y \\ \frac{dz}{d\tau} &= c \frac{a^2}{2} \end{aligned} \tag{A.20}$$

To integrate these equations of motion to obtain the trajectory let us assume that the laser field is described by

$$\mathbf{a} = a_0 \cos(\omega\tau) \hat{\mathbf{x}} \tag{A.21}$$

Integrating the x and z -components of equation (A.20) produces:

$$\begin{aligned} x(\tau) &= \frac{ca_0}{\omega} \sin(\omega\tau) \\ z(\tau) &= \frac{ca_0^2}{4} \left(\tau + \frac{1}{2\omega} \sin(2\omega\tau) \right) \end{aligned} \tag{A.22}$$

which with $\tau = t - z/c$ becomes

$$\begin{aligned} x(t) &= \frac{ca_0}{\omega} \sin(\omega t - kz) \\ z(t) &= \frac{ca_0^2}{4\omega} \left((\omega t - kz) + \frac{1}{2} \sin(2(\omega t - kz)) \right) \end{aligned} \tag{A.23}$$

A.2 The plasma disperion relation

Assuming $\mathbf{E} = E_0 e^{i(kz - \omega t)} \hat{\mathbf{x}}$ and $\mathbf{B} = B_0 e^{i(kz - \omega t)} \hat{\mathbf{y}}$ equation (A.3) can be written as

$$\epsilon_0 i \mathbf{k} \cdot \mathbf{E} = -en, \tag{A.24}$$

equation (A.2) becomes

$$i \mathbf{k} \times \mathbf{E} = -(-i\omega \mathbf{B}) = i\omega \mathbf{B}, \tag{A.25}$$

equation (A.4) can be written as

$$i \mathbf{k} \cdot \mathbf{B} = 0 \tag{A.26}$$

and equation (A.1) can be written as

$$\frac{1}{\mu_0} i \mathbf{k} \times \mathbf{B} = -ne\mathbf{v} - \epsilon_0 i\omega \mathbf{E}. \tag{A.27}$$

Using these equations and the equation of motion for a single electron in an electromagnetic field of $a_0 \ll 1$ (i.e. the $\mathbf{v} \times \mathbf{B}$ term is insignificant),

$$m_e \frac{\partial \mathbf{v}}{\partial t} = -e\mathbf{E} \tag{A.28}$$

$$\mathbf{v} = \frac{e\mathbf{E}}{m_e i\omega}. \tag{A.29}$$

Taking another curl of equation (A.25) yields: $i(\mathbf{k} \times (\mathbf{k} \times \mathbf{E})) = i\omega \mathbf{k} \times \mathbf{B}$. With $\mathbf{k} \times (\mathbf{k} \times \mathbf{E}) = \mathbf{k}(\mathbf{k} \cdot \mathbf{E}) - k^2 \mathbf{E}$, since $\mathbf{k} \cdot \mathbf{E} = 0$ for transverse waves and substituting this

into equation (A.27), the plasma dispersion relation can be found, which describes the propagation of an electromagnetic wave through an underdense plasma:

$$\omega^2 = \omega_p^2 + c^2 k^2 \quad (\text{A.30})$$

A.3 The ponderomotive force

Starting with the Lorentz force in terms of the vector and scalar potentials as given in equations (A.5) and (A.6):

$$\frac{d\mathbf{p}}{dt} = q \left[-\frac{\partial \mathbf{A}}{\partial t} - \nabla \phi + \mathbf{v} \times (\nabla \times \mathbf{A}) \right] \quad (\text{A.31})$$

and using the time derivative for the vector potential, $\frac{d}{dt} = \frac{\partial}{\partial t} + \frac{d\mathbf{r}}{dt} \cdot \frac{\partial}{\partial \mathbf{r}} = \frac{\partial}{\partial t} + \mathbf{v} \cdot \nabla$, gives:

$$\frac{d\mathbf{p}}{dt} = q \left[-\frac{d\mathbf{A}}{dt} + \mathbf{v} \cdot \nabla \mathbf{A} + \mathbf{v} \times (\nabla \times \mathbf{A}) - \nabla \phi \right]. \quad (\text{A.32})$$

Using the vector identity, $\mathbf{v} \cdot \nabla \mathbf{A} + \mathbf{v} \times (\nabla \times \mathbf{A}) = (\nabla \mathbf{A}) \cdot \mathbf{v}$, equation (A.32) can be rearranged to:

$$\frac{d}{dt}(\mathbf{p} + q\mathbf{A}) = q(\nabla \mathbf{A}) \cdot \mathbf{v} - q\nabla \phi \quad (\text{A.33})$$

With the normalisation, $\mathbf{a} = -q\mathbf{A}/mc$, $\phi \rightarrow -q\phi/mc^2$, $\mathbf{v} \rightarrow \beta = \mathbf{v}/c$ and $\mathbf{p} \rightarrow \mathbf{p}/mc$, this can be rewritten:

$$\frac{d}{dt}(\mathbf{p} - \mathbf{a}) = c\nabla \phi - (c\nabla \mathbf{a}) \cdot \mathbf{v} \quad (\text{A.34})$$

Here $\frac{d}{dt}(\mathbf{p} - \mathbf{a})$ corresponds to the rate of change of the canonical momentum, $c\nabla \phi$ is the electrostatic force due to the charge distribution and $(c\nabla \mathbf{a}) \cdot \mathbf{v}$ relates to the coupling between electromagnetic field and current density. Equation (A.34) can be written in terms of the canonical momentum $\mathbf{u} = \mathbf{p} - \mathbf{a}$ and therefore becomes:

$$\frac{d\mathbf{u}}{dt} = c\nabla \phi - (c\nabla \mathbf{a}) \cdot \frac{\mathbf{p}}{\gamma} \quad (\text{A.35})$$

$$= c\nabla \phi - (c\nabla \mathbf{a}) \cdot \frac{(\mathbf{u} + \mathbf{a})}{\gamma} \quad (\text{A.36})$$

$$\frac{d\mathbf{u}}{dt} = -\frac{1}{\gamma} c\nabla \frac{\mathbf{a}^2}{2} - c\nabla \mathbf{a} \cdot \frac{\mathbf{u}}{\gamma} + c\nabla \phi \quad (\text{A.37})$$

The first term on the RHS of equation (A.37) is an effective non-linear ponderomotive force on a single electron. To make this clearer, let us consider an interaction with an electromagnetic pulse. For a pulse with high-frequency oscillations, $\omega = |k|c$, and a spatio-temporal envelope of characteristic duration $\tau \gg 1/\omega$ and spot size $w \gg 1/|k|$, the fast and slow dynamics can be separated out. Note there are no specifications in polarisation or focusing or pulse shape of the radiation field, other than that the focusing and pulse shape must be such that the extent of the radiation field is significantly larger than the wavelength. Taking the time average of the equation over a period $\Delta T = 2\pi/\omega$ yields:

$$\left\langle \frac{d\mathbf{u}}{dt} \right\rangle = \frac{1}{\Delta T} \int_T^{T+\Delta T} dt \left[\frac{d\mathbf{u}}{dt} = -\frac{1}{\gamma} c \nabla \frac{\mathbf{a}^2}{2} - c \nabla \mathbf{a} \cdot \frac{\mathbf{u}}{\gamma} + c \nabla \phi \right] \quad (\text{A.38})$$

Using integration by parts $\int u'v dx = uv - \int uv' dx$, equation (A.38) becomes:

$$\left\langle \frac{d\mathbf{u}}{dt} \right\rangle = -\left\langle \frac{c}{\gamma} \nabla \frac{\mathbf{a}^2}{2} \right\rangle + c \nabla \langle \phi \rangle - c \nabla \langle \mathbf{a} \rangle \cdot \frac{\mathbf{u}}{\gamma} \Big|_T^{T+\Delta T} + \int_T^{T+\Delta T} c \nabla \langle \mathbf{a} \rangle \cdot \frac{\partial}{\partial t} \left(\frac{\mathbf{u}}{\gamma} \right) dt \quad (\text{A.39})$$

where $\langle \rangle$ stands for the time average. Because the envelope of the vector potential has a long scale length compared to the fast oscillations, $\langle \mathbf{a} \rangle \simeq 0$ and therefore the last two terms in equation (A.39) become zero.

The consequence is that the ponderomotive force has no polarisation dependence, it simply expels electrons from the region of high energy-density. However, if the approximation made earlier does not hold, i.e. the pulse is tightly focused and/or has few cycles, the other terms will be non-negligible and there will be a polarisation dependence to the resultant force. Since the force depends on $a^2 = q_{e,i}^2 A^2 / m_{e,i}^2 c^2$ the sign of $q_{e,i}$ does not matter, and the ponderomotive force expels both electrons and ions from regions of high field intensity.

The relativistic ponderomotive force on a single electron in a vacuum is:

$$\left\langle \frac{d\mathbf{u}}{dt} \right\rangle = \mathbf{F}'_p = -\left\langle \frac{c}{\gamma} \nabla \frac{\mathbf{a}^2}{2} \right\rangle \quad (\text{A.40})$$

The $\mathbf{F}'_p$ here is still normalised. In general, the ponderomotive force can be written as:

$$\mathbf{F}_p = -\frac{q^2}{2m \langle \gamma \rangle} \nabla \langle \mathbf{A}^2 \rangle = -\frac{mc^2}{2 \langle \gamma \rangle} \nabla \langle \mathbf{a}^2 \rangle = -\frac{mc^2}{2 \langle \gamma \rangle} \nabla (\langle \gamma \rangle^2 - 1) \quad (\text{A.41})$$

A.4 Critical power for self-focusing (P_c)

Following the derivation in [78], consider a laser beam with a radial profile

$$a(r) = a_0 \exp\left(\frac{-r^2}{\sigma_0^2}\right)$$

which has been focused to a spot size with a waist of σ_0 , just inside a region of uniform, underdense plasma, see figure A.1. In the absence of nonlinear effects, the beam will diffract with a divergence angle

$$\theta = \frac{dR}{dZ} = \frac{\sigma_0}{z_R} = \frac{2}{k\sigma_0}, \quad (\text{A.42})$$

where $z_R = \frac{\pi\sigma_0^2}{\lambda}$ is the Rayleigh length.

The relativistic refractive index before the approximation for equation (2.14) is:

$$\eta_R = \sqrt{1 - \frac{\omega_p^2}{\omega_0^2 \gamma(r)}} \quad (\text{A.43})$$

As described in section 2.2.1, when the refractive index has a peak on axis ($d\eta_R/dr < 0$), then this acts as a focusing lens. The phase velocity of the wave fronts passing through the focusing medium can be approximated using equation (A.43) and $\gamma = \sqrt{1 + a^2(r)/2} \simeq 1 + a^2(r)/4$:

$$\frac{v_{ph}(r)}{c} = \frac{1}{\eta_R} \simeq 1 + \frac{\omega_p^2}{2\omega_0^2} \left(1 - \frac{a^2(r)}{4}\right) \quad (\text{A.44})$$

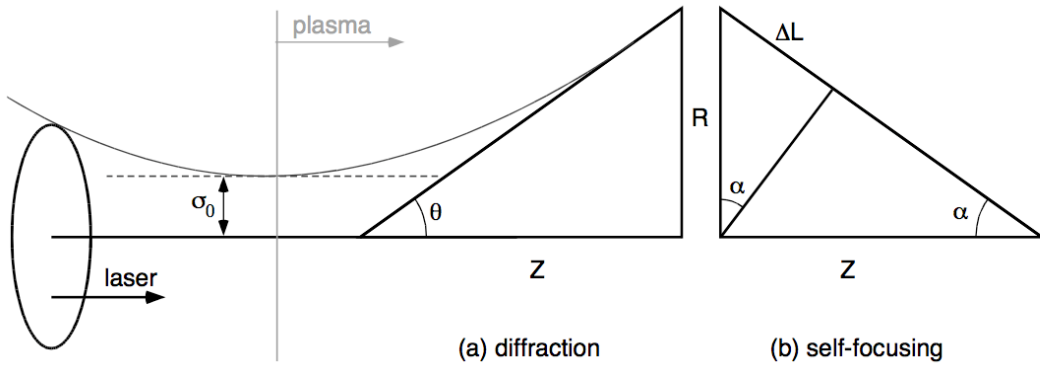


Figure A.1: The effect of (a) diffraction and (b) self-focusing during the interaction between a high-intensity laser and an underdense plasma.

Looking across the beam profile, see figures A.1 and 2.4, the phase fronts will travel more slowly at the centre than at the edge, giving a velocity difference:

$$\frac{\Delta v_{ph}(r)}{c} = \frac{\omega_p^2}{8\omega_0^2} a_0^2 \exp(-2r^2/\sigma_0^2)$$

This curvature in the phase front causes the rays to bend by an amount dependent on their relative path difference. The maximum path difference is:

$$\Delta L = |\Delta v_{ph}|_{\max} t = \left| \frac{\Delta v_{ph}}{c} \right|_{\max} Z = \alpha R.$$

Therefore the maximum focusing angle of the beam is given by:

$$\alpha^2 = \frac{\omega_p^2 a_0^2}{8\omega_0^2} \quad (\text{A.45})$$

Beam expansion due to diffraction can therefore be cancelled by self-focusing when $\theta = \alpha$. Setting equations (A.42) and (A.45) equal gives, using $k = \omega_0/c$:

$$a_0^2 \frac{\omega_p^2}{8\omega_0^2} = \frac{4c^2}{\omega_0^2 \sigma_0^2} \quad (\text{A.46})$$

The expression for the beam power P_L in watts is:

$$\begin{aligned} P_L &= \int_S I \, d^2r \\ &= \pi c \epsilon_0 \int_0^\infty E^2(r) r \, dr \end{aligned} \quad (\text{A.47})$$

With $E(r) = m_e \omega_0 c a(r)/e$ this becomes:

$$P_L = \frac{\pi c^3 \epsilon_0 m_e^2 \omega_0^2 a_0^2 \sigma_0^2}{4e^2} \quad (\text{A.48})$$

Rearranging equation (A.46) for $a_0^2 \sigma_0^2$, and putting the result in equation (A.48) one obtains:

$$P_L \geq \frac{8\pi \epsilon_0 m_e^2 c^5 \omega_0^2}{e^2 \omega_p^2} \quad (\text{A.49})$$

This corresponds to the critical power:

$$P_c = \frac{8\pi \epsilon_0 m_e^2 c^5 \omega_0^2}{e^2 \omega_p^2} = 17 \left(\frac{n_c}{n_e} \right) \text{ [GW]} \quad (\text{A.50})$$

A.5 The plasma frequency

The equation of motion for an electron is:

$$\frac{d\mathbf{p}}{dt} = q[\mathbf{E} + \mathbf{v} \times \mathbf{B}] \quad (\text{A.51})$$

Using the fact that $\frac{d}{dt} = \frac{\partial}{\partial t} + (\mathbf{v} \cdot \nabla)$ this becomes:

$$\frac{\partial \mathbf{v}}{\partial t} + (\mathbf{v} \cdot \nabla) \cdot \mathbf{v} = \frac{q}{m_e} [\mathbf{E} + \mathbf{v} \times \mathbf{B}] \quad (\text{A.52})$$

To derive the plasma frequency (following Chen [72]) we make the following assumptions: $kT = 0$, no thermal motion, the ions are heavy, and therefore do not move, no B -field, motion only in x -direction, and the plasma is infinitely large. With $\nabla = \hat{\mathbf{x}} \frac{\partial}{\partial x}$, $\mathbf{E} = E\hat{\mathbf{x}}$, $\nabla \times \mathbf{E} = 0$ and $\mathbf{E} = -\nabla\phi$, the equation of motion (see equation (A.52)) yields

$$m_e n_e \left(\frac{\partial \mathbf{v}_e}{\partial t} + (\mathbf{v}_e \cdot \nabla) \mathbf{v}_e \right) = -e \mathbf{E} n_e \quad (\text{A.53})$$

and the continuity equation becomes:

$$\frac{\partial n_e}{\partial t} + \nabla(n_e \mathbf{v}_e) = 0 \quad (\text{A.54})$$

The Poisson Equation is given by:

$$\epsilon_0 \nabla E = e(n_i - n_e) \quad (\text{A.55})$$

With $n_e = n_0 + n_1$, $\mathbf{v}_e = \mathbf{v}_0 + \mathbf{v}_1$ and $\mathbf{E}_e = \mathbf{E}_0 + \mathbf{E}_1$, because

$$\begin{aligned} \nabla n_0 &= 0 = \mathbf{v}_0 = \mathbf{E}_0 \\ \frac{\partial n_0}{\partial t} &= 0 = \frac{\partial \mathbf{v}_0}{\partial t} = \frac{\partial \mathbf{E}_0}{\partial t} \end{aligned}$$

by putting this the above into equation (A.53), we obtain:

$$m_e \left(\frac{\partial \mathbf{v}_1}{\partial t} + (\mathbf{v}_1 \cdot \nabla) \mathbf{v}_1 \right) = -e \mathbf{E}_1 \quad (\text{A.56})$$

The term $(\mathbf{v}_1 \cdot \nabla) \mathbf{v}_1 \rightarrow 0$ since it is seen to be quadratic in an amplitude quantity, and we shall linearise by neglecting it.

Therefore the preceding equation becomes:

$$m_e \frac{\partial \mathbf{v}_1}{\partial t} = -e \mathbf{E}_1 \quad (\text{A.57})$$

The continuity equation becomes:

$$\frac{\partial n_1}{\partial t} + \nabla \cdot ((n_0 + n_1)\mathbf{v}_1) = 0 \quad (\text{A.58})$$

$$\frac{\partial n_1}{\partial t} + n_0 \nabla \cdot \mathbf{v}_1 = 0 \quad (\text{A.59})$$

$$\epsilon_0 \nabla E_1 = e(n_{i0} + n_{i1} - n_{e0} - n_{e1}) = -en_1 \quad (\text{A.60})$$

because $n_{i0} \rightarrow n_{e0}$. The oscillating quantities are assumed to behave sinusoidally:

$$\begin{aligned} \mathbf{v}_1 &= v_1 e^{i(kx - \omega t)} \hat{\mathbf{x}} \\ n_1 &= n_1 e^{i(kx - \omega t)} \\ \mathbf{E}_1 &= E_1 e^{i(kx - \omega t)} \hat{\mathbf{x}} \end{aligned} \quad (\text{A.61})$$

Equations (A.56), (A.58) and (A.59) become:

$$m_e i \omega v_1 = e E_1 \quad (\text{A.62})$$

$$-n_1 i \omega + i n_0 k v_1 = 0 \quad (\text{A.63})$$

$$\epsilon_0 i k E_1 = -e n_1 \quad (\text{A.64})$$

Eliminating n_1 and E_1 we have for equation (A.62):

$$-i m_e \omega v_1 = -e \frac{-e}{i k \epsilon_0} \frac{-n_0 i k v_1}{-i \omega} = -i \frac{n_0 e^2}{\epsilon_0 \omega} v_1 \quad (\text{A.65})$$

$$\Rightarrow \omega_p = \sqrt{\frac{n_0 e^2}{\epsilon_0 m_e}} \quad (\text{A.66})$$

A.6 The Rayleigh length

The Rayleigh length is the length over which the laser stays in focus. For a gaussian pulse it is defined as the distance from focus where the intensity in the area of the pulse cross section drops by a factor of 2. $z_{RG} = \frac{\pi w_0^2}{\lambda}$ for a Gaussian beam where w_0 is the waist, which in turn is defined where the intensity drops radially by a factor of $\frac{1}{e^2}$.

M^2 -factor is a parameter for quantifying the beam quality of a laser [112].

$$M^2 = \frac{w_{0R} \theta_{0R}}{w_{0G} \theta_{0G}} = \frac{\pi}{\lambda} \arctan(0.5f) w_{0R} \simeq \frac{\pi}{2\lambda} \frac{w_{0R}}{f}$$

where f is the f-number and θ_{0R} and θ_{0G} are the far field divergence angles of the real beam and the Gaussian beam, respectively. For a non Gaussian beam, $z_R = \frac{z_{RG}}{M^2}$, $M^2 = 1$ corresponds to a gaussian beam. A beam is usually M^2 times diffraction limited. With $\text{FWHM} \approx w_0$, one obtains $z_R = 2fd_{\text{FWHM}}$.

A.7 Betatron resonance

Using [23, 25, 142] and considering a plane electromagnetic wave, propagation in the z -direction, polarised in the x -direction. Additionally imposing a static electric field $E_x^s = \kappa_E x$ and a magnetic field $B_y^s = -\kappa_B x$. The transverse electric field of the laser is:

$$E_x^L = E_0 \cos(\omega t - kz)$$

with $\omega/k = v_{ph}$

$$E_x^L = E_0 \cos(\omega t - c/v_{ph})$$

$$\mathbf{E}^L = -\frac{\partial \mathbf{A}}{\partial t} = E_0 \cos(k_L z - \omega_L t) \hat{\mathbf{x}} \quad \text{with } E_0 = \omega_L A_0$$

$$\mathbf{B}^L = \nabla \times \mathbf{A} = B_0 \cos(k_L z - \omega_L t) \hat{\mathbf{y}} \quad \text{with } B_0 = k_L A_0 = \eta/c E_0$$

Putting $k_L = \omega_L/v_{ph} = \omega_L \eta/c$ into $B_0 = k_L A_0$ with $\omega_L A_0 = E_0$ gives,

$$\Rightarrow B_0 = E_0/v_{ph}$$

$$E_x^L = E_0 \cos(\omega t - \omega z/v_{ph}) \text{ and } B_y^L = E_x^L/v_{ph}$$

$$\frac{d\mathbf{p}}{dt} = -[(\mathbf{E}^L + \mathbf{E}^s) + (\mathbf{v} \times \mathbf{B}^L + \mathbf{v} \times \mathbf{B}^s)] \quad (\text{A.67})$$

$$\begin{pmatrix} v_x \\ v_y \\ v_z \end{pmatrix} \times \begin{pmatrix} B_x \\ B_y \\ B_z \end{pmatrix} = \begin{pmatrix} v_y B_z - v_z B_y \\ v_z B_x - v_x B_z \\ v_x B_y - v_y B_x \end{pmatrix} = \begin{pmatrix} -v_z B_y \\ 0 \\ v_x B_y \end{pmatrix} \quad (\text{A.68})$$

The z component of (A.67) is:

$$\frac{dp_z}{dt} = -[E_z + (v \times B^L + v \times B^s)]$$

With the z component of equation (A.68) and $E_z=0$ this becomes:

$$\frac{dp_z}{dt} = -0 - (v_x B_y^L + v_x B_y^s)$$

$$\frac{dp_z}{dt} = -\frac{v_x}{v_{ph}} E^L - v_x B^s \quad (\text{A.69})$$

The x component of (A.67) with the x component of equation (A.68) becomes:

$$\frac{dp_x}{dt} = -[E_x^L + E_x^s + (-v_z B_y^L - v_z B_y^s)]$$

B^L can be substituted by E^L/v_{ph} .

$$\frac{dp_x}{dt} = -\left(1 - \frac{v_z}{v_{ph}}\right) E^L - E^s + v_z B^s \quad (\text{A.70})$$

Equation (A.69) and (A.70) can be put together for the energy conservation: with $\gamma = \sqrt{1 + p_x^2 + p_z^2} \rightarrow \frac{d\gamma}{dt} = \frac{1}{2\gamma} \left(\frac{dp_z^2}{dt} + \frac{dp_x^2}{dt} \right) = \frac{1}{2\gamma} \left(2p_z \frac{dp_z}{dt} + 2p_x \frac{dp_x}{dt} \right)$ and using $p = \gamma v$ this then becomes:

$$\frac{d\gamma}{dt} = -v_x (E^L + E^s) \quad (\text{A.71})$$

The transverse momentum can be written as $p_x = \gamma v_x$. Therefore

$$\frac{dp_x}{dt} = \frac{d(\gamma v_x)}{dt} = v_x \frac{d\gamma}{dt} + \gamma \frac{dv_x}{dt},$$

equating this with equation (A.70), substituting equation (A.71) for $\frac{d\gamma}{dt}$ as well as $\mathbf{E}^s = \kappa_E x \hat{\mathbf{x}}$ and $\mathbf{B}^s = -\kappa_B x \hat{\mathbf{x}}$ leads to

$$\frac{d^2 x}{dt^2} + (\kappa_E + v_z \kappa_B) x / \gamma = \left[\left(\frac{dx}{dt} \right)^2 - \left(1 - \frac{v_z}{v_{ph}} \right) \right] \frac{E^L}{\gamma} + \left(\frac{dx}{dt} \right)^2 \frac{\kappa_E x}{\gamma} \quad (\text{A.72})$$

Here $\omega_\beta^2 = (\kappa_E + v_z \kappa_B) / \gamma$ is the betatron frequency and γ is the Lorentz factor of the accelerated electron. Equation (A.72) describes a driven oscillator. A resonance will occur when the frequency of the driving term on the right hand side and the betatron frequency are the same, i.e. when:

$$\omega_\beta = 1 - \frac{v_z}{v_{ph}} \quad (\text{A.73})$$

Equation (A.73) is normalised to ω_0 . This occurs when the laser frequency, as witnessed by an electron moving at a velocity v_z , is the same as the oscillation or bounce frequency of the electron in the channel. The resonance condition states that when an electron makes one oscillation, the electromagnetic wave, which propagates

with $v_{ph} > c$, overtakes it exactly by one period. Thus the electron and electromagnetic wave remain in phase and efficient energy gain is possible. When the electrons are accelerated beyond resonance, they start to decelerate again. Using the dispersion relation $\omega_0^2 = \omega_p^2 + c^2 k^2$ and $v_{ph} = \frac{\omega_0}{k}$, the phase velocity of the laser can be written as:

$$\frac{v_{ph}}{c} = \left(1 - \frac{\omega_p^2}{\langle\gamma\rangle\omega_0^2}\right)^{-1/2} \simeq 1 + \frac{1}{2\langle\gamma\rangle} \frac{\omega_p^2}{\omega_0^2} \quad (\text{A.74})$$

where γ is the Lorentz factor of the background electrons which affects the laser propagation. For a relativistic electron with $v_z \rightarrow c$ in an underdense plasma resonance occurs when:[†]

$$\omega_\beta \simeq \frac{1}{2\langle\gamma\rangle} \frac{n_e}{n_c} \quad (\text{A.75})$$

To calculate the electron energy at which the resonance occurs, γ_e , we use a general expression for the betatron frequency,

$$\omega_\beta^2 = \frac{\nabla \mathbf{F}}{\gamma_e} \quad (\text{A.76})$$

where $\mathbf{F}$ is the focusing force (due to the combination of magnetic and electric fields). In the simple case above the focusing force was linear but in general this will not be the case. Since $\langle\gamma\rangle = \sqrt{1 + a_0^2/2}$ for linearly polarised light[‡], the resonance condition can be expressed in terms of the background plasma density and laser intensity as:

$$\gamma_e \simeq 4 \frac{n_c^2}{n_e^2} \left(1 + \frac{a_0^2}{2}\right) \nabla \mathbf{F} \quad (\text{A.77})$$

Assuming $\nabla \mathbf{F} \propto a_0$ (see chapter 4.3.1) for a given density yields, $\gamma_e \propto a_0 + a_0^3/2$.

[†]From equation (A.74) to equation (A.75): $\omega_\beta = \left(1 - \frac{v_z}{v_{ph}}\right)$, with $v_z \rightarrow c$, and $\frac{v_{ph}}{c} \approx \left(1 + \frac{1}{2\langle\gamma\rangle} \frac{\omega_p^2}{\omega_0^2}\right)$, ω_β therefore becomes:

$$\omega_\beta = \left(1 - \frac{c}{v_{ph}}\right) \approx 1 - \left(1 + \frac{1}{2\langle\gamma\rangle} \frac{\omega_p^2}{\omega_0^2}\right)^{-1} = \frac{1 + \frac{1}{2\langle\gamma\rangle} \frac{\omega_p^2}{\omega_0^2} - 1}{1 + \frac{1}{2\langle\gamma\rangle} \frac{\omega_p^2}{\omega_0^2}}$$

The denominator of the right hand side goes to 1. With γ very large and n_e very small for under-dense plasmas.

[‡] $\langle\gamma\rangle = \sqrt{1 + \langle\mathbf{a}^2\rangle}$ for linear polarisation $\rightarrow \langle\gamma\rangle = \sqrt{1 + \frac{a_0^2}{2}}$

A.8 Capacitor approximation for vacuum heating

One model frequently used to describe the vacuum heating mechanism is the capacitor approximation [78]. For this, the magnetic field of the laser is neglected (i.e. the $\mathbf{v} \times \mathbf{B}$ term is neglected). It also assumes that the laser has an electric field component that is perpendicular to the target surface, E_d , and therefore pulls the electrons back and forth across their equilibrium positions. A schematic of this model is shown in figure A.2. The wave, with electric field E_L is obliquely incident

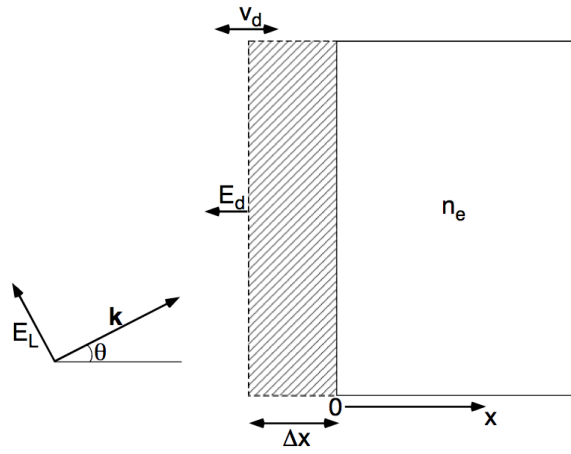


Figure A.2: Capacitor model of the Brunel heating mechanism.

onto an initially smooth surface. Considering that a reflected wave will be present, which is for the moment assumed to have the same amplitude as the incoming wave, the driving electric field is given by:

$$E_d = 2E_L \sin \theta \quad (\text{A.78})$$

This field pulls out electrons, and displaces them by a distance Δx from their original position, leading to a *surface number density* of the sheet of $\Sigma = n_e \Delta x$, as for capacitor plates. The electric field created between $x = -\Delta x$ and $x = 0$ can be found by:

$$\Delta E = -\frac{1}{\epsilon_0} e \int_{-\Delta x}^{x=0} n_e dx = \frac{1}{\epsilon_0} e n_e \Delta x = \frac{1}{\epsilon_0} e \Sigma \quad (\text{A.79})$$

Setting $E_d = \Delta E$ and solving for Σ gives:

$$\Sigma = \frac{2\epsilon_0 E_L \sin \theta}{e} \quad (\text{A.80})$$

The returning charge sheet will have acquired a velocity $v_d \simeq 2v_{os}\sin\theta$, where v_{os} is the electron quiver velocity in the laser field and $a_0 = v_{os}/c$. Using the assumption that all electrons enter the solid target and no longer feel the electric field of the laser, the average energy density absorbed per laser cycle is:

$$\begin{aligned} I_a &= \frac{\Sigma m_e v_d^2}{\tau} \frac{1}{2} \\ &= \frac{\epsilon_0}{4\pi} \frac{e}{m_e \omega_0} E_d^3 \end{aligned} \quad (\text{A.81})$$

For the conversion in equation (A.81), $\tau = 2\pi/\omega_0$ and $v_{os} = \frac{eE_L}{m_e \omega_0}$ are used. The incoming laser intensity is given by $I_x = I_{in}\cos\theta = \frac{\epsilon_0 c}{2} E_L^2 \cos\theta = \frac{c}{8\pi} E_L^2 \cos\theta$. Comparing this to equation (A.81), gives the fractional absorption rate:

$$f_{bh} = \frac{I_a}{I_x} = \frac{4}{\pi} a_0 \frac{\sin^3\theta}{\cos\theta} \quad (\text{A.82})$$

Taking a closer look at equation (A.82) shows that it predicts an (unlimited) increase of absorption with larger angles of incidence and higher laser irradiance ($I\lambda^2 \propto a_0^2$). To correct for this, the first step is to assume imperfect reflectivity, and therefore replacing equation (A.78) with:

$$E_d' = \left[1 + \sqrt{1 - f_{bh}}\right] E_L \sin\theta \quad (\text{A.83})$$

In the following $b = [1 + \sqrt{1 - f_{bh}}]$ will be substituted. Note that this also influences $v_d' \simeq b v_{os}\sin\theta$ and equation (A.80) becomes $\Sigma' = \frac{b \epsilon_0 E_L \sin\theta}{e}$. The second correction takes the relativistic effects for intensities $\geq 10^{18} \text{ Wcm}^{-2}$ into account, and leads to a replacement of the kinetic energy with $U_k = (\gamma - 1)m_e c^2$, where $\gamma = \sqrt{1 + v_d'^2/c^2}$. This leads to equation (A.81) taking the form:

$$\begin{aligned} I_a' &= \frac{\Sigma'}{\tau} (\gamma - 1) m_e c^2 \\ &\simeq \frac{b^3 \epsilon_0}{4\pi} \frac{e}{m_e \omega_0} E_L^3 \sin^3\theta \end{aligned} \quad (\text{A.84})$$

Then the absorption rate becomes:

$$f_{bh}' = \frac{I_a'}{I_x} = \frac{b^3}{2\pi} a_0 \frac{\sin^3\theta}{\cos\theta} = \frac{b}{\pi a_0} \left[\sqrt{1 + b^2 a_0^2 \sin^2\theta} - 1 \right] \frac{\sin\theta}{\cos\theta} \quad (\text{A.85})$$

Assuming that $a_0 \gg 1$, then $ba_0\sin\theta \gg 1$ in this strongly relativistic case, and

$$f_{bh}'' = \frac{b^2 \sin^2\theta}{\pi \cos\theta}. \quad (\text{A.86})$$

With $\alpha'' = \frac{\sin^2\theta}{\cos\theta}$ and from $b = [1 + \sqrt{1 - f_{bh}}] \rightarrow f_{bh} = 2b - b^2$, equation (A.86) can be reduced to a simple quadratic for b , and one finds $b = 2(\frac{\alpha''}{\pi} + 1)^{-1}$, and:

$$f_{bh}^{rel} = \frac{4\pi\alpha''}{(\pi + \alpha'')^2} \quad (\text{A.87})$$

A.9 Derivation of the electric sheath field

Following [141], we assume that N_e electrons are accelerated by the laser and are confined in an electron bunch of length $L=c\tau$. At the rear side of the target the electrons are transversely spread over a circular area with radius

$$r_B = r_L + d \tan\theta,$$

where r_L is the radius of the laser spot, d is the target thickness and θ is the half-angle of the electrons traveling through the target.

We further assume an exponential electron energy distribution:

$$\frac{dN}{d\mathcal{E}} = \frac{N_e}{k_B T_e} \exp\left(-\frac{\mathcal{E}}{k_B T_e}\right)$$

When electrons cross the solid-vacuum boundary, they induce a positive surface charge Qe on the conducting rear surface, leading to a surface charge density $Qe/(\pi r_B^2)$ located at $z = 0$, where z is the electron propagation axis.

Solving the Poisson equation for such a charge density distribution, the potential on the z -axis is given by:

$$-e\Phi(\xi) = \mathcal{E}_\infty s(\xi) \quad (\text{A.88})$$

with $\xi = z/r_B$ and $s(\xi) = 1 + \xi - \sqrt{1 + \xi^2}$. Only a few electrons with the energies exceeding $\mathcal{E}_\infty = Qe^2/(2\pi\epsilon_0 r_B)$ can escape the rear surface potential. The low energetic electrons re-enter the foil. The point $\hat{\xi} = \hat{z}/r_B$ where electrons with the mean energy $k_B T_e$ turn around is defined by $\mathcal{E}_\infty s(\hat{\xi}) \approx \mathcal{E}_\infty \hat{\xi} = k_B T_e$, for $\hat{\xi} \ll 1$.

Using the definition for $\mathcal{E}_\infty$ and the electron density $n_{Q0} = Q/(\pi r_B^2 \hat{z})$ directly at the surface, one derives:

$$\hat{z} = \sqrt{\frac{2\epsilon_0 k_B T_e}{n_{Q0} e^2}} \equiv \lambda_{D\text{hot}}, \quad (\text{A.89})$$

where $\lambda_{D\text{hot}}$ is called the hot Debye length.

From equation (A.88) together with $E = -d\Phi/dz$, we obtain:

$$E_z(\xi) = \frac{k_B T_e}{e \lambda_{D\text{hot}}} \left(1 - \frac{\xi}{\sqrt{1 + \xi^2}} \right) \quad (\text{A.90})$$

A.10 Derivation for laser peak intensity

For a calculation of the peak intensity, I_0 , one has to consider both the temporal profile of the pulse as well as the spatial profile of the focus. For a gaussian profile $f(x) = e^{-b \frac{x^2}{x_0^2}}$, the factor b can be found for the HWHM using $f(x_0) = \frac{1}{2}$: $\Rightarrow b = \ln 2$. Assuming a Gaussian profile for both the temporal and spatial profile, the intensity can be written as:

$$I(r, t) = I_0 e^{-\ln 2 \frac{r^2}{r_0^2}} e^{-\ln 2 \frac{4t^2}{\tau^2}} \quad (\text{A.91})$$

where I_0 is the peak intensity, r is the distance from the centre of focus, r_0 is the HWHM of the focal spot, t is the time and τ is the FWHM of the pulse duration. From this, by integrating over the area, we can find the total power:

$$P(t) = \int I(r, t) dA = \int_0^{2\pi} \int_0^\infty I(r, t) r dr d\theta \quad (\text{A.92})$$

with $\int_0^\infty e^{-b'x^2} x dx = \frac{1}{2b'}$ this gives:

$$P(t) = \frac{\pi r_0^2}{\ln 2} I_0 e^{-\ln 2 \frac{4t^2}{\tau^2}} \quad (\text{A.93})$$

By integrating once more over time, we find the total energy in the pulse:

$$\mathcal{E} = \int_{-\infty}^\infty P(t) dt \quad (\text{A.94})$$

With $\int_{-\infty}^\infty e^{-b'x^2} dx = \sqrt{\frac{\pi}{b'}}$ we get:

$$\mathcal{E} = I_0 \frac{r_0^2 \tau}{2} \left(\frac{\pi}{\ln 2} \right)^{\frac{3}{2}} \quad (\text{A.95})$$

r_0 can be substituted by $r_0 = d/2$ where d is the FWHM. Equation (A.95) then becomes:

$$\mathcal{E} = I_0 \frac{d^2 \tau}{8} \left(\frac{\pi}{\ln 2} \right)^{\frac{3}{2}} \quad (\text{A.96})$$

Solving for I_0 gives:

$$I_0 = \frac{8 \mathcal{E}_{laser}}{\tau d^2} \left(\frac{\ln 2}{\pi} \right)^{\frac{3}{2}} = 0.829 * \frac{\mathcal{E}_{laser}}{\tau d^2} \quad (\text{A.97})$$

Appendix B

Tracking code and analytical solution

B.1 Analytical solution

```
Bmax = 1.1; % Tesla
rB = 11e-3; % radius
detector_plane = 250e-3
E_min = 10; % minimum electron energy in MeV
E_max = 500; % maximum electron energy in MeV
N_electrons = 30; % number of different energy electrons to track
%constants
me = 9.11e-31; % electron mass in kg
m_MeV = 0.511; % electron mass in MeV
e = 1.6e-19;
c = 2.998e8;
%assigning energies to the electrons
E_e_MeV = (1:N_electrons);
o=(E_max - E_min)/(N_electrons-1);
E_e_MeV(1)=E_min;
for i=2:N_electrons
    E_e_MeV(i)=E_e_MeV(i-1)+o;
end
gamma =(E_e_MeV./m_MeV)+ 1;
r_analytical=zeros(length(E_e_MeV),1);
angle=zeros(length(E_e_MeV),1);
for i=1:length(E_e_MeV)
    angle(i)=2.*atan((Bmax*rB*e)./(abs(sqrt((gamma(i).^2)-1)*me*c))); %atan result in radians!
    r_analytical(i)=detector_plane.*tand(angle(i)); %tand result in deg
end
figure; plot(E_e_MeV,r_analytical(:,1)); xlabel('Electron energy in MeV'); ylabel('deflection in m'); title('Analytical')
```

B.2 Test tracking code

```
Bmax = 1.1 ; %magnetic field in Tesla
detector_plane = 300; %distance from centre of magnet to detector plane in mm
plane = 250; %real distance from centre of magnet to detector plane in mm real since onlt the fastest reach detector_plane
%plane is the distance from centre of magnet to the
```

```

%plane where you want to monitor the electrons
detector.height = 30; % height of detectors
target = 1000; % distance from target to magnet centre in mm
collimation = 0.01; % collimation half angle (radians)
E.min = 10; % minimum electron energy
E.max = 500; % maximum electron energy
N.electrons = 30; % number of different energy electrons to track
dBy = 40; %y-dimension of magnet in mm
dBx = 40; %x-dimension of magnet in mm (direction of propagation)
%%% constants
m.MeV = 0.511; %electron mass in MeV
me = 9.11e-31; %electron mass in Kg
c = 2.998e8;
e = 1.6e-19;
%%%
B.test_map=zeros(dBy,dBx);
r = 11;
M = 1000;

x_mag = linspace(0,(dBx-1), dBx).*1e-3;
y_mag = linspace(0,(dBy-1), dBy).*1e-3;
x_mag = x_mag - (dBx/2)*1e-3;
y_mag = y_mag - (dBy/2)*1e-3;
[x_mag, y_mag]= meshgrid(x_mag, y_mag);
x0 = round(length(B.test_map(1,:))/2);
y0 = round(length(B.test_map(:,1))/2);
%%set up grid from magnet centre to detector plane
x.min = -detector.plane*1e-3;
x.max = detector.plane*1e-3;
y.min = -detector.height*1e-3;
y.max = detector.height*1e-3;
dx = 1e-3; %spacing (1mm)
dy = 1e-3;
x = [x.min:dx:x.max];
y = [y.min:dy:y.max];
[x.array,y.array] = meshgrid(x,y); %"plaid" arrays

for q=1:length(B.test_map(1,:))
    for v=1:length(B.test_map(:,1))
        t=sqrt((q-x0).^2+(v-y0).^2);
        if r>t
            B.test_map(v,q)=1;
        elseif r==t
            B.test_map(v,q)=1;
        elseif r<t
            B.test_map(v,q)=0;
        end
    end
end

B.test = griddata(x_mag,y_mag,B.test_map,x.array,y.array);

%%get rid of NaN in B.test
B.testM=isnan(B.test);
for i=1:length(y)
    for j=1:length(x)
        if B.testM(i,j)>0
            B.test(i,j)=0;
        end
    end
end

%make every value 1 or 0

```

```

for q=1:length(B.test(1,:))
    for v=1:length(B.test(:,1))
        if B.test(v,q)>0
            if B.test(v,q)<1
                B.test(v,q)=1;
            end
        end
    end
end

E_e_MeV = (1:N_electrons);
o=(E_max - E_min)/(N_electrons-1);
E_e_MeV(1) = E_min;
for i=2:N_electrons
    E_e_MeV(i)=E_e_MeV(i-1)+o;
end
gamma_e = (E_e_MeV./m.MeV) + 1;
p_e_0= sqrt(gamma_e.^2 - 1) * m_e * c; % momentum in SI units
dtheta = collimation;
t = 0.0;
dt = 1e-11;
Nsteps = fix((abs(max(max(x_array)))+abs(min(min(x_array))))/2.4832e8/dt);

p_e = zeros(N_electrons,2,Nsteps); %time history of momentum
%p_e(i,j,k) i is index for which electron, j = 1 is x coordinate, j = 2 is y coordinate and k is t coordinate
r_e = zeros(N_electrons,2,Nsteps); %time history of positions
v_e = zeros(N_electrons,2,Nsteps);
B = zeros(N_electrons,Nsteps);

for theta=0 %~dtheta:dtheta:+dtheta
    %set up the initial positions and momenta
    yi = (target - detector_plane)*tan(theta)*1e-3;
    xi = -detector_plane*1e-3;
    pi_x = p_e_0 .*cos(theta);
    pi_y = p_e_0 .*sin(theta);
    p_e(:,1,1)=pi_x;
    p_e(:,2,1)=pi_y;
    r_e(:,1,1)=xi;
    r_e(:,2,1)=yi;
    gammanew=gamma_e(:);
    B(:,1)=interp2(x_array,y_array, B_test,xi,yi);
    v_e(:,1,1) = p_e(:,1,1)./gammanew(:)./me;
    v_e(:,2,1) = p_e(:,2,1)./gammanew(:)./me;
    p_e(:,1,2)=p_e(:,1,1)-e*dt.*(v_e(:,2,1).*B(:,1));
    p_e(:,2,2)=p_e(:,2,1)+e*dt.*(v_e(:,1,1).*B(:,1));
    v_e(:,1,2) = p_e(:,1,2)./gammanew(:)./me;
    v_e(:,2,2) = p_e(:,2,2)./gammanew(:)./me;
    r_e(:,1,2)=r_e(:,1,1)+dt*v_e(:,1,1);
    r_e(:,2,2)=r_e(:,2,1)+dt*v_e(:,2,1);
    B(:,2)=interp2(x_array,y_array, B_test, r_e(:,1,2),r_e(:,2,2));

    for i = 3:Nsteps

        p_e(:,1,i)=p_e(:,1,i-2)-e*2*dt.*(v_e(:,2,i-1).*B(:,i-1));
        p_e(:,2,i)=p_e(:,2,i-2)+e*2*dt.*(v_e(:,1,i-1).*B(:,i-1));
        v_e(:,1,i) = p_e(:,1,i)./gammanew(:)./me;
        v_e(:,2,i) = p_e(:,2,i)./gammanew(:)./me;
        r_e(:,1,i)= r_e(:,1,i-1)+dt*v_e(:,1,i-1);
        r_e(:,2,i)= r_e(:,2,i-1)+dt*v_e(:,2,i-1);

        if r_e(:,2,i)<max(max(y_array))
            if r_e(:,2,i)<abs(min(min(y_array)))
                B(:,i) = interp2(x_array,y_array, B_test, r_e(:,1,i),r_e(:,2,i));
            end
        end
    end
end

```

```

    %get Bfield value at current electron position by interpolation
    elseif r_e(:,2,i)==abs(min(min(y_array)))
        B(:,i) = interp2(x_array,y_array, B_test, r_e(:,1,i),r_e(:,2,i));
    elseif r_e(:,2,i)>abs(min(min(y_array)))
        B(:,i)=0;
    end
    elseif r_e(:,2,i)==max(max(y_array))
        if r_e(:,2,i)<abs(min(min(y_array)))
            B(:,i) = interp2(x_array,y_array, B_test, r_e(:,1,i),r_e(:,2,i));
        elseif r_e(:,2,i)==abs(min(min(y_array)))
            B(:,i) = interp2(x_array,y_array, B_test, r_e(:,1,i),r_e(:,2,i));
        elseif r_e(:,2,i)>abs(min(min(y_array)))
            B(:,i) = 0;
        end
    elseif r_e(:,2,i)>max(max(y_array))
        B(:,i) = 0;
    end

    t = t +dt;
end

for i=1:N_electrons
    r_ex1 = r_e(i,1,:);
    r_ey1 = r_e(i,2,:);
    r_ex1 = r_ex1(:);
    r_ey1 = r_ey1(:);
    plot(r_ex1,r_ey1)
    hold on
end

hold off
maybedeflection = zeros(N_electrons,1,Nsteps);
deflection = zeros(N_electrons,2);

%get y positions at the detector plane and not the max stepnumber
for j=1:N_electrons
    for i=1:Nsteps
        if round(r_e(j,1,i).*1e2)==plane*1e-1
            maybedeflection(j,1,i)=r_e(j,1,i);
        end
    end
    delta=abs(plane*1e-3- maybedeflection(j,1,:));
    for k=2:Nsteps
        if min(delta(1,1,k))==min(delta(1,1,:))
            pos=k;
        end
    end
    deflection(j,:,1)=r_e(j,:,pos);
end

figure
plot(E_e.MeV,deflection(:,2,1))
xlabel('Electron energy in MeV')
ylabel('deflection in m')
title('Btest')

%%%FIT
xdata=E_e.MeV(:);
if theta==dtheta
    ydata=deflection(:,2,1)+collimation;
else
    ydata=deflection(:,2,1);
end
M1=1;

```

```

M2=0.5;
[estimates, model] = fitcurvedemo(xdata,ydata);
figure
plot(xdata, ydata, 's')
hold on
[sse, FittedCurve] = model(estimates);
plot(xdata, FittedCurve, 'r')
xlabel('Energy MeV')
ylabel('deflection in meter')
title(['a=', num2str(estimates(1)), ' ', 'b=', num2str(estimates(2))]);
legend('data', ['fit using ', 'power law'])
hold off
end

```

For the different electron spectrometers and configurations, there are different tracking codes, but all are based on the code described in this section. They mostly vary in magnetic field inputs, and how these files including the magnetic field profiles are read into the program. The electron beam divergence is taken into account. It can be different for different experiments. For all tracking codes one gets three values for a and b , one for the on axis case, and two for $+d\phi$ and $-d\phi$.

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Publications

The following publications present work discussed in this thesis:

1. S. KNEIP, **S.R. Nagel**, C. BELLEI, N. BOURGEOIS, A.E. DANGOR, A. GOPAL, R. HEATHCOTE, S.P.D. MANGLES, J.R. MARQUES, A. MAKSIMCHUK, P.M. NILSON, K.T. PHUOC, S. REED, M. TZOUFRAS, F.S. TSUNG, L. WILLINGALE, W.B. MORI, A. ROUSSE, K. KRUSHELNICK AND Z. NAJMUDIN, *Observation of synchrotron radiation from electrons accelerated in a petawatt-laser-generated plasma cavity*, Physical Review Letters **100**(10) 105006 (2008)
2. L. WILLINGALE, **S.R. Nagel**, A.G.R. THOMAS, C. BELLEI, R.J. CLARKE, A.E. DANGOR, R. HEATHCOTE, M.C. KALUZA, C.KAMPERIDIS, S.KNEIP, K. KRUSHELNICK, N. LOPES, S.P.D. MANGLES, W. NAZAROV, P.M. NILSON AND Z. NAJMUDIN, *Characterization of high-intensity laser propagation in the relativistic transparent regime through measurements of energetic proton beams*, Physical Review Letters **102**(12) 125002 (2009)

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S.R. Nagel, S.P.D. MANGLES, S. KNEIP, C. BELLEI, R.J. CLARKE, A. E. DANGOR, A. GOPAL, R. HEATHCOTE, K.L. LANCASTER, A. MAKSIMSHUK, S.A. REED, M. TATARAKIS, L. WILLINGALE, K. KRUSHELNICK AND Z. NAJMUDIN *Aspects of direct laser acceleration in ultra-intense laser plasma interactions*. in preparation

The following publications present work was performed during the PhD qualification period, but is not presented in this thesis:

3. C. BELLEI, **S.R. Nagel**, S. KAR, A. HENIG, S. KNEIP, C. PALMER, A. SAEVERT, L. WILLINGALE, D. CARROLL, B. DROMEY, J.S. GREEN, K. MARKEY, P.T. SIMPSON, R.J. CLARKE, D. NEELY, M. KALUZA, S.P.D. MANGLES, P. MCKENNA, P.A. NORREYS, J. SCHREIBER, M. ZEPF, K. KRUSHELNICK, AND Z.

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4. S. KNEIP, **S.R. Nagel**, S. MARTINS, C. BELLEI, O. CHEKLOV, R.J. CLARKE, A.E. DANGOR, N. DELERUE, E.J. DIVALL, G. DOUCAS, K. ERTTEL, F. FIUZA, R. FONSECA, P. FOSTER, S.J. HAWKES, C.J. HOOKER, K. KRUSHELNICK, W.B. MORI, C. PALMER, K. TA PHUOC, P. RAJEEV, J. SCHREIBER, M.J.V. STREETER, D. URNER, J. VIEIRA, L.O. SILVA, S.P.D. MANGLES, AND Z. NAJMUDIN, *Near-GeV acceleration of electrons by a non-linear plasma wave driven by a self-guided laser pulse*, submitted (2009)
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