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Department of Physics

Relativistic Plasma Photonics: Polarisation Control and Higher-Order Mode Generation in Laser-Solid Interactions



by

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in Physics

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Abstract

This thesis reports on experimental and numerical investigations of relativistic plasma photonic effects arising from the interaction of high-intensity ($> 10^{18}$ W/cm²) laser pulses with solid targets. The work focuses on two themes: the production of frequency-doubled light with higher-order spatial mode structures and the control of the polarisation state of light transmitted through ultra-thin (tens of nanometer) targets.

The first investigation explores how aperture targets can be used to generate intense second-harmonic light with structured spatial profiles. Two- and three-dimensional EPOCH simulations identify the target thickness and aperture diameter that maximise conversion efficiency, with the resulting mode shape shown to depend on the laser polarisation. Further improvements were achieved using tapered apertures matched to the convergence angle of the incoming beam.

The second study examines practical challenges in producing $2\omega_L$ light with aperture targets. Simulations using EPOCH and FLASH are used to explore the effects of plasma forming within the aperture and small misalignments of the laser beam arising from spatial jitter. Both reduce conversion efficiency, but jitter is shown to have more influence, and possible strategies to minimise it are discussed.

The third investigation is based on experimental campaign with the Gemini laser, testing the generation of $2\omega_L$ light from aperture targets. Measurements of transmitted ω_L and generated $2\omega_L$ are compared with simulations, confirming that beam jitter and focal-spot quality strongly influence conversion. The results demonstrate the importance of stabilising pointing and preserving high-quality Gaussian beam profiles in future experiments.

The final study explores changes to the polarisation of light transmitted through ultra-thin (tens of nanometer) targets irradiated by intense pulses, using the Vulcan laser. A secondary, lower-intensity pulse was used to pre-expand the targets, altering transmission. An unexpected relationship between transmission and polarisation was

observed. Comparisons with analytical models suggest that spatio-temporal variations in the transmitted light, combined with diagnostic sub-sampling, explain this effect. These findings highlight the importance of capturing the full transmitted beam when diagnosing polarisation in ultra-thin target experiments.

Overall, the research demonstrates how target design, laser stability and diagnostic precision together determine the effectiveness of relativistic plasma photonic control.

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Role of Author

The results presented in this thesis are the work of the author over several experimental campaigns, performed as part of a team of researchers led by Prof. Paul McKenna. The author had a primary role in the execution of the experimental campaigns, data analysis, numerical simulations and interpretation of results.

Chapter 4: The author performed the 2D EPOCH simulations under the supervision of Dr. Martin King (University of Strathclyde) and Dr. Ross Gray (University of Strathclyde). The 3D EPOCH simulations were performed collaboratively with Dr King. The author led the analysis of all simulation data, including the development of the FFT function used to filter the simulation data for a selected frequency. The author developed the mathematical model describing the efficiency of laser pulse light to frequency-doubled light in collaboration with Dr. King. Use of the BISHOP code was supported by Dr. Gray, also the original developer of this code.

Chapter 5: The author was primarily responsible for running the FLASH and EPOCH simulations, with guidance from Dr. King on the latter when required. The author led the analysis of all simulation data, with support from Dr. Gray and Dr. King.

Chapter 6: The author actively contributed to the planning and execution of the experiment and took a primary role in the analysis of all experimental and simulation data, with support from Dr. Gray and Dr. King. FLASH and EPOCH simulations were primarily conducted by the author, with Dr. King providing advice on the EPOCH simulations.

Chapter 7: The author contributed to the execution of the experiment and assisted in constructing the polarimetry diagnostic, designed by Dr. Timothy Frazer (University of Strathclyde). The author took a leading role in the analysis of all experimental and simulation data, with support from Dr. Gray and Dr. King. The 2D EPOCH simulations performed by the author, with Dr. King providing advice when required.

The 3D EPOCH simulations were performed by Dr. King.

Publications

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Chapter 1

Introduction

The research presented in this thesis sits at the intersection of two complementary areas of study. The first, *high-power laser-solid interactions*, examines how intense laser pulses interact with solid density matter and the extreme energy densities involved ionise the material to produce a plasma, a state of matter with highly tunable optical properties. The second, *relativistic plasma photonics*, seeks to exploit these properties by using plasma as a tuneable optical medium to generate and control high-intensity light. Unlike conventional solid-state optics, plasma has no intrinsic damage threshold, making it exceptionally well suited to manipulating light at intensities that would otherwise destroy traditional materials. This thesis investigates how plasma produced by intense laser-solid interactions can be harnessed to develop new techniques for manipulating and generating high-intensity light.

This introductory chapter first provides a brief, non-comprehensive overview of the development of intense laser pulses. It then discusses the fields of laser-solid interactions and relativistic plasma photonics, including representative examples of relevant processes, thereby establishing the context for the investigations presented in the later chapters.

1.1 History of intense laser pulse development

The history of the laser begins with the discovery of the phenomenon of stimulated emission, the release of a photon from an excited atom induced by interaction with an external photon. First proposed by Albert Einstein in 1917 [1] and subsequently observed experimentally by Rudolf Walther Ladenburg in 1928 [2], no immediate ap-

plications were apparent. However, in the mid 20th century, research into exploiting stimulated emission as a means of generating light intensified, with one prominent researcher in this field, Gordon Gould, introducing the term ‘laser’– Light Amplification by Stimulated Emission of Radiation [3]. This work led to the demonstration of the first operational laser by Theodore H. Maiman in 1960 [4]. The distinctive properties of laser light, such as its monochromaticity and coherence, motivated the continued research into both the refinement and enhancement of laser technology. As a result, lasers now find application across a variety of fields such as communications, industry, and medicine, from fibre optics [5] to precision machining [6] and surgical techniques [7].

Of particular relevance to the work presented in this study is the continual drive toward achieving higher laser power and intensity. This is motivated by the desire of achieving higher energy densities that not only advance fundamental physics but also lay the groundwork for novel applications in science and technology. Key developments enabling this progress included the advent of Q-switching [8] and mode-locking [9], which facilitated the production of short-duration, high-energy laser pulses. Over the subsequent decades, steady improvements in laser technology led to progressively higher pulse intensities. Figure 1.1 shows the increase in laser peak-power over time, with key technologies and facilities indicated. A key point in this development is the surpassing of the ionisation threshold of matter ($\sim 10^{10}$ W/cm²). This enabled the generation of plasma by such pulses, and thereby the development of laser-matter interactions as an area of research. As this figure demonstrates, the laser peak-power plateaued at the $\sim$ GW level for several decades, corresponding to pulse intensities around $\sim 10^{17}$ W/cm². This limitation was imposed by the damage thresholds of the solid-state optics used within such lasers, preventing further amplification. This barrier was overcome with the development of chirped pulse amplification (CPA), first demonstrated in 1985 [10], which enabled the amplification of ultrashort pulses without exceeding optical damage thresholds.

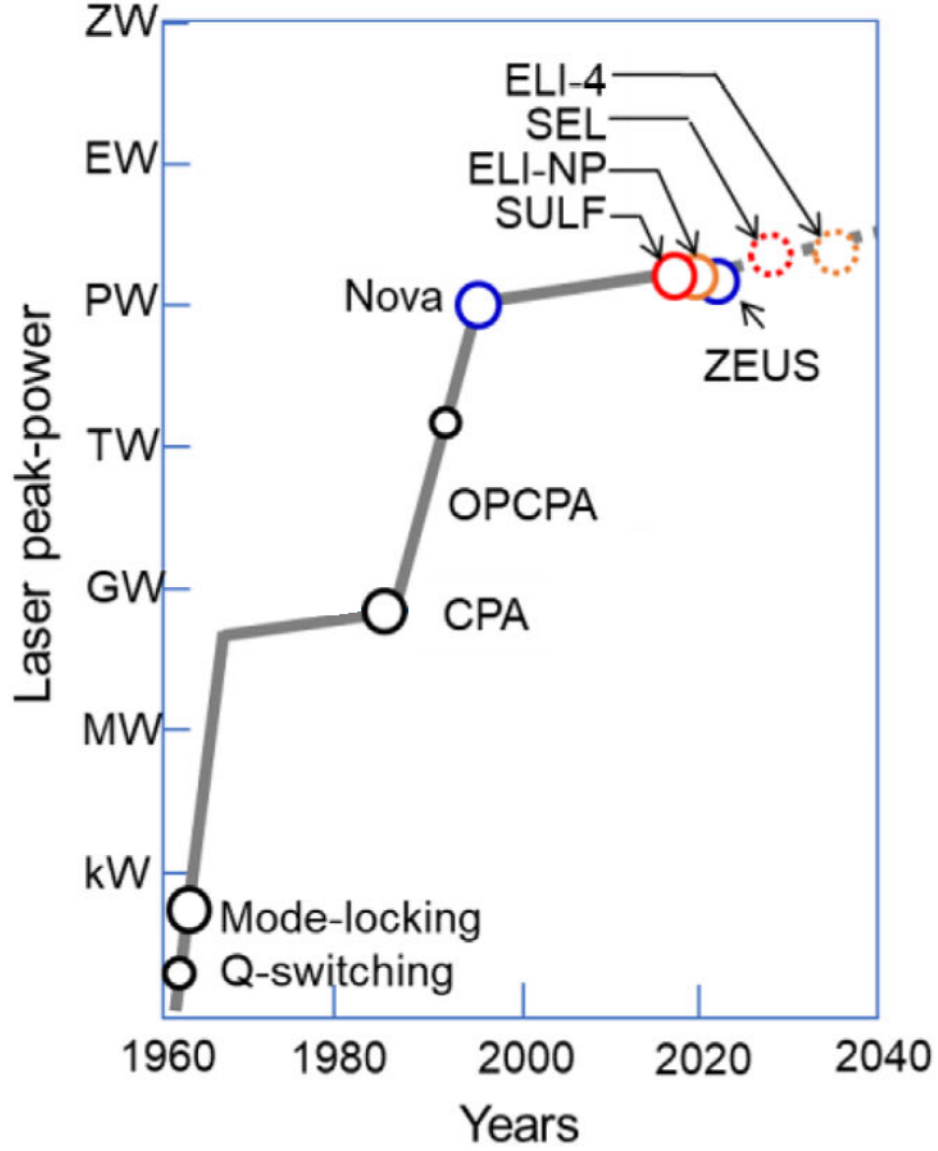


Figure 1.1: Increases in laser peak-power since 1960. Key laser technologies and facilities are indicated. Adapted from [11].

CPA made possible the generation of laser pulses with intensities exceeding 1×10^{18} W/cm². At a laser wavelength of $\sim 1 \mu\text{m}$, such an intensity is capable of accelerating electrons to relativistic energies in one laser period and are therefore termed ‘relativistically intense’. This allowed for access to new regimes of high-intensity laser physics. In the subsequent decades, there has been a sustained effort to further increase laser peak-power. One key advance was the development of an additional amplification scheme known as optical parametric chirped pulse amplification (OPCPA) [12]. First demonstrated in A. Dubietis et al. [13] in 1992, it was later adopted in practical high power laser facilities at the Central Laser Facility (CLF) in 1997 [14], the location of the laser

systems employed for the experimental campaigns reported on in chapters 6 and 7 of this thesis. This progress has paved the way for the current generation of high-power laser facilities, such as the 10 PW Shanghai Superintense Ultrashort Laser Facility (SULF) in China [15], the 10 PW Extreme Light Infrastructure - Nuclear Physics (ELI-NP) [16] located in the Romania, the 10 PW Apollon system located in France [17] and the 3 PW Zettawatt-Equivalent Ultrashort pulse laser System (ZEUS) located in Michigan, USA [18] being capable of routinely generating pulses with peak intensities in the range of $\sim 10^{22}$ W/cm². In addition, next generation facilities, such as the 20 PW ‘Vulcan 2020’ project [19] and the 1 PW Extreme Photonics Application Center (EPAC) [20], both located at the CLF, are under construction. As of writing, the highest recorded peak pulse intensity is 1×10^{23} W/cm², achieved using the CoReLS laser facility in South Korea [21]. In addition to the increase in peak laser pulse power, there has been a push toward increasing the repetition rate of high-power laser systems. This allows for larger volumes of data and, consequently, more robust statistical analysis, enabling physical effects that are rare or obscured by noise to be revealed. For example, the SCAPA facility at the University of Strathclyde [22] possesses a 350 TW laser system capable of generation pulses at 5 Hz (although typically operated at 1 Hz), while the Central Laser Facility’s most recent installation, EPAC, is designed to deliver pulses at a repetition rate of 10 Hz. These advances, while remarkable, also revealed the intrinsic limitations of solid-state optics - thus giving rise to the emergence of plasma-based photonic solutions capable of withstanding relativistic intensities.

1.2 Relativistic plasma photonics

The peak pulse intensities now achievable at current generation high-power laser facilities far exceed the damage threshold of solid-state optics. Consequently, the manipulation of such light presents a significant challenge, potentially limiting the exploration of new research frontiers within high-intensity laser science and the applications of such high-intensity pulses. Given that this is a fundamental limitation of solid-state optics, research into alternative means of controlling high-intensity light is ongoing, with plasma emerging as a promising candidate. Unlike conventional materials, plasma can withstand extremely high energy densities, and its optical properties can be modified on the ultra-fast timescale of a high-power laser pulse, allowing for dynamic manip-

ulation of intense light. This concept, known as plasma optics or plasma photonics, is an emerging field at the intersection of plasma physics and photonics, focused on exploiting plasmas optical properties to generate and control light [23].

Plasma, as a state of matter, can exist over a wide range of densities and temperatures, and plasma photonic techniques are found across this spectrum. Many involve relatively low-temperature plasmas, for example, a discharge lamp could be considered a plasma photonic technique [24]. However, such cases lie outside the scope of the research presented here, which focuses on relativistic plasma photonics. A plasma is said to be relativistic when its electrons are accelerated to velocities approaching the speed of light, where effects such as relativistic mass increase strongly influence plasma dynamics and its optical response. This, combined with the absence of a traditional damage threshold, makes plasma a promising medium for controlling high-intensity light.

A particularly effective means of generating relativistic plasma is through laser-solid interactions. A relativistically intense drive pulse incident on a solid target induces ionisation through mechanisms such as multiphoton [25] and barrier suppression ionisation [26]. This process produces a plasma suitable for studying relativistic effects. This occurs during the initial stages of the interaction, leaving a substantial portion of the laser pulse that accelerates the plasma electrons to near-light speeds through its strong electromagnetic fields. Laser-solid interactions under these conditions are particularly well suited for investigating relativistic plasma photonic techniques, as the nature of the plasmas produced are inherently suitable for this purpose. Depending on interaction parameters such as target geometry and laser intensity, plasma photonic effects such as modification of the drive pulse properties and secondary radiation generation can occur. The following subsections briefly survey key areas where plasma has been proposed as a photonic medium: amplification (1.2.1), conditioning (1.2.2), polarisation control (1.2.3), and the generation of structured light (1.2.4). The latter is a central focus of this thesis.

1.2.1 Pulse amplification

A key objective of current research into high-power laser technology is the push towards higher peak powers, and therefore greater intensities [27]. When this is achieved through increased amplification of the pulse, the solid-state optics used within the CPA

process (section 3.2.1) must be increased in size such that the energy density of the pulse does not exceed the damage threshold of such optics. However, the cost of fabricating optics on this scale makes this approach prohibitively expensive and in addition, technological challenges (e.g. in the size to which crystals can be grown) are also limiting factors. This has motivated research into plasma-based amplification schemes that are not constrained by the damage thresholds of conventional optics. One of the key properties that makes plasma suitable for this task is that, unlike solid materials, it has no damage threshold in the conventional sense.

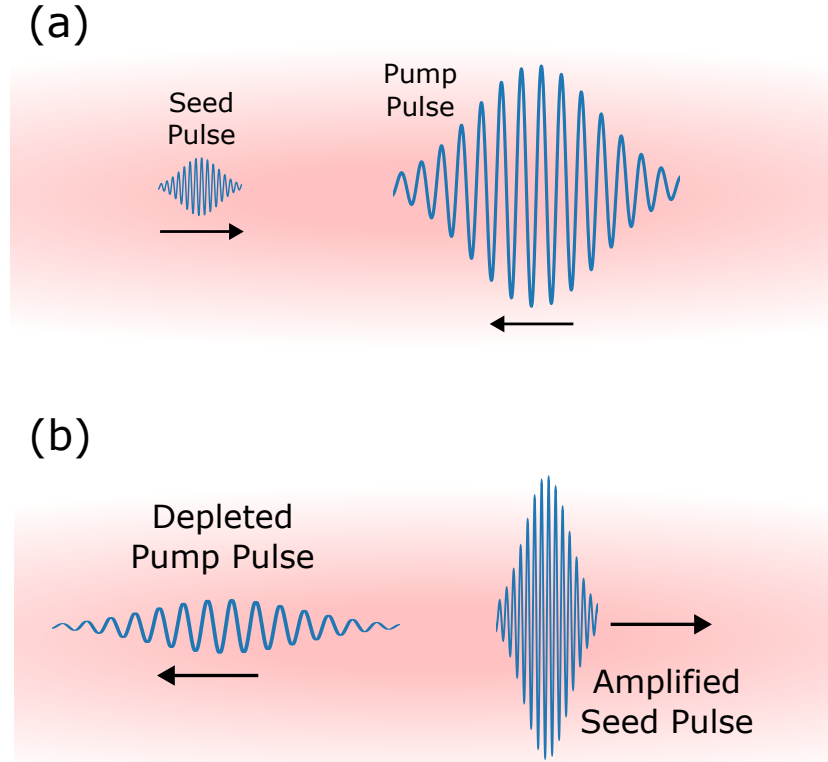


Figure 1.2: Illustration of laser amplification via stimulated Raman scattering. The seed and pump pulse are shown counter-propagating through a plasma (a) before and (b) after interacting.

One such technique is amplification via stimulated Raman scattering (SRS) [28], often termed Raman amplification. In this process, two laser pulses are counter-propagated through a plasma. The pulse being amplified, referred to as the seed pulse, typically has a short duration, in the tens to hundreds of femtosecond range. The second pulse, acting as the energy source and referred to as the pump pulse, is much longer in duration, ranging from tens of picoseconds to nanosecond. Raman scattering is the inelastic scattering of photons by matter. When occurring within a plasma, this event

results in the generation of both a seed photon, with a frequency matching the seed pulse, and a plasma wave. The plasma wave then mediates the transfer of energy from the pump to seed pulse. Figure 1.2 shows an illustration of this process.

This process was first proposed in Malkin *et al.* [29], in which a theoretical model for the transfer of energy between the pump and seed pulse is presented. Trines *et al.* [30] demonstrates numerically the possibility of reaching pulse powers in the range of hundreds of petawatts. Experimental investigations have since demonstrated the feasibility of this form of amplification. Ping *et al.* [31] demonstrated a $95\times$ amplification of a seed pulse via Raman amplification in underdense plasmas. Marques *et al.* [32] demonstrated amplification of sub-picosecond pulses above the Joule level with an energy transfer efficiency between the two pulses of around 20%. J. Ren *et al.* [33] demonstrated the amplification of a seed pulse from $1.3 \times 10^{14} \text{ W/cm}^2$ to $\sim 4 \times 10^{16} \text{ W/cm}^2$ by employing a double-pass SRS amplification scheme.

Despite the many successful experimental demonstrations of Raman amplification, the promise of reaching energies equalling or exceeding those achievable using CPA has yet to be demonstrated. Studies have been conducted into the conditions required to make this feasible, such as Trines *et al.* [34] who provide criteria to improve the efficiency of SRS amplification. The authors suggest that previous experimental studies have not employed the necessary interaction distances, in addition to an incorrect focus on the ‘gain factor’, the ratio of the amplified seed pulse energy to its initial value. Instead, the focus should be on ensuring that the triple product of the seed pulse amplitude, duration, and the Raman scattering coupling constant exceeds a critical threshold.

1.2.2 Pulse conditioning

The laser pulses employed in laser-solid interactions are typically preceded by background light that, when brought to focus, is of sufficient intensity to induce ionisation of the target prior to the arrival of the main pulse. This can negatively impact interaction dynamics, so techniques are often implemented to suppress its intensity. A plasma photonic component widely employed for this task is the planar plasma mirror (PPM) [35]. This consists of a planar piece of dielectric glass that is transparent to the laser wavelength. When positioned before the laser focus, the PPM ensures that the preceding light, having an intensity below the ionisation threshold, is largely transmitted, while the main pulse induces ionisation of its surface. The generated plasma reflects

the remainder of the main pulse, which is subsequently employed to drive the primary laser-solid interaction. In this way, the preceding light is reduced in intensity by several orders of magnitude. As key component in the experiments reported on in this thesis, a more detailed discussion is provided in section 3.2.2.

While the PPM is the most widely employed, alternative geometries have been investigated that seek to provide an additional degree of control over the temporal and spatial characteristics of the reflected pulse. One such example is the so-called ‘Focusing plasma mirror’ (FPM) that aims to simultaneously focus the pulse while enhancing its contrast. Nakatsutsumi *et al.* [36] demonstrated a factor 5 decrease in focal spot size through the use of a F/0.4 ellipsoidal plasma mirror in combination with an OAP, compared to the sole use of an OAP. Wilson *et al.* [37, 38] demonstrated a factor of 2.5 reduction in focal spot size through the use of F/1 ellipsoidal plasma mirrors in combination with an F/3 OAP, again in comparison to the sole use of an OAP. However, this study also showed that the focal spot quality is highly sensitive to misalignment with the FPM, arising both from inaccuracies in FPM placement relative to the OAP focus and from shifts in the OAP focal position caused by wavefront aberrations. In both cases, intensity enhancement was gauged by the maximum proton energy obtained via TNSA [39] when these pulses were incident on foil targets. As this is well-characterised acceleration mechanism [40], this enabled inference of the pulse intensity.

Plasma photonic techniques for the compression of pulses have also been widely studied. A crucial stage of the CPA process is the compression of the pulse after amplification of the pulse has occurred, covered in further detail in section 3.2.1. Similar to the amplification of pulses discussed in this chapter, this is typically achieved using solid-state optical components which are susceptible to damage from the high-energy density pulse. Again, increasing the size of the components responsible for compression is prohibitively expensive, motivating the search for a plasma-based alternative.

One such approach is the use of a plasma Bragg grating (PBG) [41]. In conventional photonics, a Bragg grating is formed by creating periodic variations in refractive index, whereas in the plasma analogue, this variation is created through plasma density modulation as the refractive index of a plasma is proportional to its electron density. An example of the formation of such a grating is through the interference of two or more laser pulses [42, 43]. The resulting interference pattern gives rise to periodic vari-

ations in laser intensity. When incident upon a material, varying rates of ionisation are induced by the higher and lower intensity regions, leading to corresponding variations in plasma density. The feasibility of pulse compression using PBG's has been demonstrated via PIC simulation [44, 45, 41], and experimental evidence of the generation of plasma gratings and subsequent diffraction has been demonstrated [46].

Plasma photonic techniques for the purpose of focusing light have been investigated. Again, motivated by the damage thresholds of conventional optical components. One such technique, termed 'plasma-lensing', exploits a phenomenon known as self-focusing. This is a non-linear optical effect in which the refractive index of a medium becomes intensity-dependent, causing a high-intensity laser pulse to behave as if it were propagating through a lens, thereby concentrating the light into a smaller spatial region. In solid-state materials, this effect is generally undesirable, as it can lead to peak intensities exceeding the damage threshold of the material. In a plasma, self-focusing can be harnessed to compress the pulse spatially, reduce its focal spot size, and potentially achieving higher intensities than are accessible with solid-state optics.

Figure 1.3 shows an illustration of the focusing of a pulse via a plasma lens, used in combination with a solid-state optic. Zeng *et al.* [47] demonstrate theoretically and numerically the use of a plasma lens in combination with a solid-state focusing system to alter the focal spot size. In this configuration, the plasma lens acts in a similar way to an eyepiece on a telescope. Edwards *et al.* [48] demonstrate how through the use of overlapping pulses within a gas jet, a so-called 'holographic' plasma lens can be generated. Through theory and numerical simulation, it is demonstrated that this technique has the potential to achieve increases in pulse intensity of several orders of magnitude. Molavi Choobini and Ghaffari-Oskooei [49] investigate, via theory and numerical simulation, the self-focusing behaviour of Gaussian pulses within a magnetised plasma. A key finding is a difference in the minimum focal spot size achieved depending on the 'handedness' of the circularly polarised light being focused. Palastro *et al.* [50] provide a theoretical and numerical study of the chromatic and non-linear phase aberrations that occur as high-intensity pulses are focused within a plasma, demonstrating their deleterious effects on the maximum achievable intensity using this technique. Gupta *et al.* [51] investigate the use of an upward plasma density ramp in an effort to prevent the expansion of the pulse after self-focusing within a plasma, allowing for the propagation of the self-focused pulse before expansion.

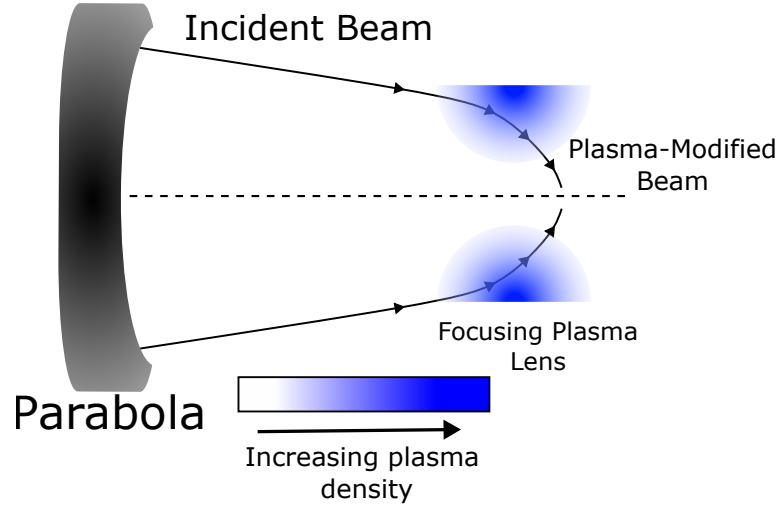


Figure 1.3: Illustration of combined solid-state parabolic mirror and plasma lens systems capable of focusing a pulse.

1.2.3 Polarisation control

Control over the polarisation state of light is often achieved using optical components such as wave plates or polarisers. Similar to the case of an amplifier gain medium, such components consisting of solid-state materials are damaged upon interaction with high-intensity light above the damage threshold. As such, research into plasma photonic techniques for controlling the polarisation state of light is largely driven by the medium's lack of damage threshold.

The control of the optical polarisation state of high-intensity light by inducing birefringence within a plasma has been investigated. A medium is said to be birefringent when its refractive index is dependent on both the polarisation and propagation direction of light passing through it. When induced in a plasma, this medium can behave as a polariser. Stark *et al.* [52] investigated, both theoretically and numerically, the generation of birefringence within a plasma arising from temperature anisotropy. The interaction between a linearly polarised pulse and a plasma results in the preferential heating of electrons along the pulse polarisation axis, inducing this anisotropy. As the plasma refractive index is proportional to the electron temperature, this leads to different refractive indices along orthogonal directions.

Another plasma photonic technique for controlling the polarisation state of light makes use of the Faraday effect, the rotation of the polarisation state of light as it

propagates through a medium subject to an applied magnetic field. While more often used as a diagnostic to infer the strengths of magnetic fields within plasmas, its ability to alter the polarisation state of light has been investigated as a potential plasma photonic technique. As with other plasma photonic techniques, a key advantage is the ability to alter the polarisation state of high-intensity pulses that would otherwise exceed the damage thresholds of conventionally used solid-state optics. Weng *et al.* [53] demonstrate theoretically and numerically the splitting of a high-intensity linearly polarised pulse into circularly polarised light after propagating through a magnetised plasma.

The polarisation of a pulse can also be controlled through wave mixing, in a somewhat similar setup to SRS amplification. Through numerical and theoretical investigation, Michel *et al.* [54] reported on the ability to control the polarisation state of a ‘probe’ pulse as it overlaps with a ‘pump’ pulse within a plasma. Their overlap leads to a modulation of the plasma refractive index, with control of the probe pulse polarisation state achieved by modifying the amplitude of the pump pulse. This technique has since been demonstrated experimentally in a paper by Turnbull *et al.* [55] in which elliptically polarised light was converted to circularly polarised after being co-propagated with a second pulse through an underdense plasma.

1.2.4 Structured light and higher order modes

Structured light refers to laser beams whose phase, amplitude, or polarisation are deliberately shaped to produce complex spatial or spatio-temporal structures. Examples include higher order transverse spatial modes, electromagnetic field distributions in the transverse plane that differ from the fundamental Gaussian mode, or so-called ‘vortex beams’ carrying orbital angular momentum (OAM). The generation of light with such properties can be achieved with solid-state components such as spatial light modulators or diffractive optics. However, the use of such components with high-intensity light is prevented by their damage threshold, motivating the exploration of plasma-photonic methods for generating structured light.

A key paper by our group and directly relevant to the research presented here, Duff *et al.* [56] reports on the use of ultra-thin (tens of nanometers) targets for the generation of frequency-doubled light possessing higher order transverse spatial modes. The initial interaction between a relativistically intense drive laser pulse possessing a

Gaussian transverse spatial intensity distribution and an ultra-thin target leads to the formation of a so-called ‘relativistic plasma aperture’. This phenomenon, introduced by our group in Gonzalez-Izquierdo *et al.* [57], involves a region of the target, at the most intense part of the laser focus, becoming relativistically transparent to the incident laser pulse, enabled by the relativistic mass increase experienced by the plasma electrons, due to a phenomenon known as relativistic self-induced transparency (RSIT) [58, 59, 60]. The remaining portion of the pulse propagates through the self-induced relativistic plasma aperture, interacting with its edges and thereby accelerating electron bunches [61]. The acceleration and deceleration of these bunches in electrostatic fields at the rear of the target results in the generation of both fundamental and frequency doubled light, with the latter possessing higher order spatial modes. An additional plasma-photonic effect that results from the formation of the plasma aperture is an observed shift in the polarisation state of the light at the rear of the target. This is attributed to the superposition of the transmitted pulse and the fundamental frequency generated light.

Yi [62] investigated, both numerically and theoretically, the generation of high-harmonic light possessing OAM from the interaction between relativistically intense pulses and aperture targets with thickness on the order of hundreds of nanometers. This is attributed to the oscillation of plasma that forms at the aperture edges, acting as a so-called “relativistic oscillating window”. The generation of high-harmonic light possessing OAM from targets with geometries differing from a simple aperture have also been studied. Trines *et al.* [63] presents a theoretical and numerical study demonstrating how tailored structured targets can generate high-harmonic light with independently tunable frequency and OAM in both reflection and transmission.

In summary, the emergence of relativistic plasma photonics provides a promising route to overcome the material limitations inherent to traditional optics at extreme intensities. By investigating how plasma produced in laser-solid interactions can generate structured and polarisation-controlled high-intensity light, this thesis seeks to advance both our fundamental understanding of light-plasma coupling and the practical development of next-generation optical technologies.

1.3 Thesis Outline

This study presents research comprising both numerical and experimental investigations into relativistic plasma photonics building directly from the work present in Duff *et al* [56] . Two central themes are addressed: (1) The generation of high-intensity, frequency-doubled light with higher order transverse spatial modes produced via laser-solid interactions. (2) The use of laser-solid interactions as a means of controlling the polarisation state of high-intensity light. The structure of this thesis is as follows:

- **Chapter 2** - An exploration of the physics underlying laser-plasma interactions required to understand the work present in the later chapters of this thesis.
- **Chapter 3** - An overview of the experimental techniques and numerical simulations used to achieve the results present in this thesis.
- **Chapter 4** - Numerical investigation of higher order modes of intense second harmonic light generated from a plasma aperture
- **Chapter 5** - Numerical simulations to guide the experimental investigation of $2\omega_L$ light generated from aperture targets
- **Chapter 6** - Experimental investigation of the generation of frequency-doubled light from pre-formed aperture targets
- **Chapter 7** - Investigation of the spatio-temporal variations in the polarisation of light transmitted through relativistically induced plasma-apertures
- **Chapter 8** - Conclusions & future work

Chapter 2

Fundamentals of laser-solid interactions

To fully appreciate the research presented in the later chapters of this thesis, it is essential to understand the fundamental physics underpinning the field of laser-plasma interactions (LPI). While LPI physics encompasses a wide range of subfields with overlaps into other research domains, this thesis focuses specifically on mode generation and polarisation control achieved through plasma-photonic techniques. Accordingly, this chapter addresses only the current understanding of the physics necessary to comprehend these topics.

This chapter begins with a discussion of the fundamental properties of high-power laser pulses used in typical laser-solid interactions studied in this field, followed by an exploration of the ionisation processes that occur when these pulses interact with matter, leading to plasma formation. The subsequent sections explore the behaviour and propagation of light within the resulting plasma, and the mechanisms by which energy from the laser pulse is absorbed into the plasma.

2.1 High power lasers

2.1.1 Pulse parameters

Before delving into the specific physics common to all LPI's, it is necessary to understand the characteristics of the laser pulses generated by high-power laser systems utilised in this field which are classified as short-duration and high-intensity. While these terms can carry different meanings across various fields of physics, in this con-

text they correspond to the following specific ranges. With respect to the duration of the pulse, ‘short’ refers to timescales between tens of femtoseconds and single-digit picoseconds. A pulse’s duration can be understood by examining the way in which its electric field strength varies over time, referred to as its temporal intensity profile. This generally consists of a sinusoidally varying electric field encapsulated within a Gaussian envelope. In the field of LPI, the conventional measure of pulse duration is the full-width-at-half-maximum (FWHM) of this envelope, often represented by the symbol t_L . An example of an idealised temporal intensity profile of a pulse with $t_L = 40$ fs is shown in figure 2.1(a), a duration typical of the pulses produced by the Gemini laser system, the details of which will be covered in a later chapter (3.3.1).

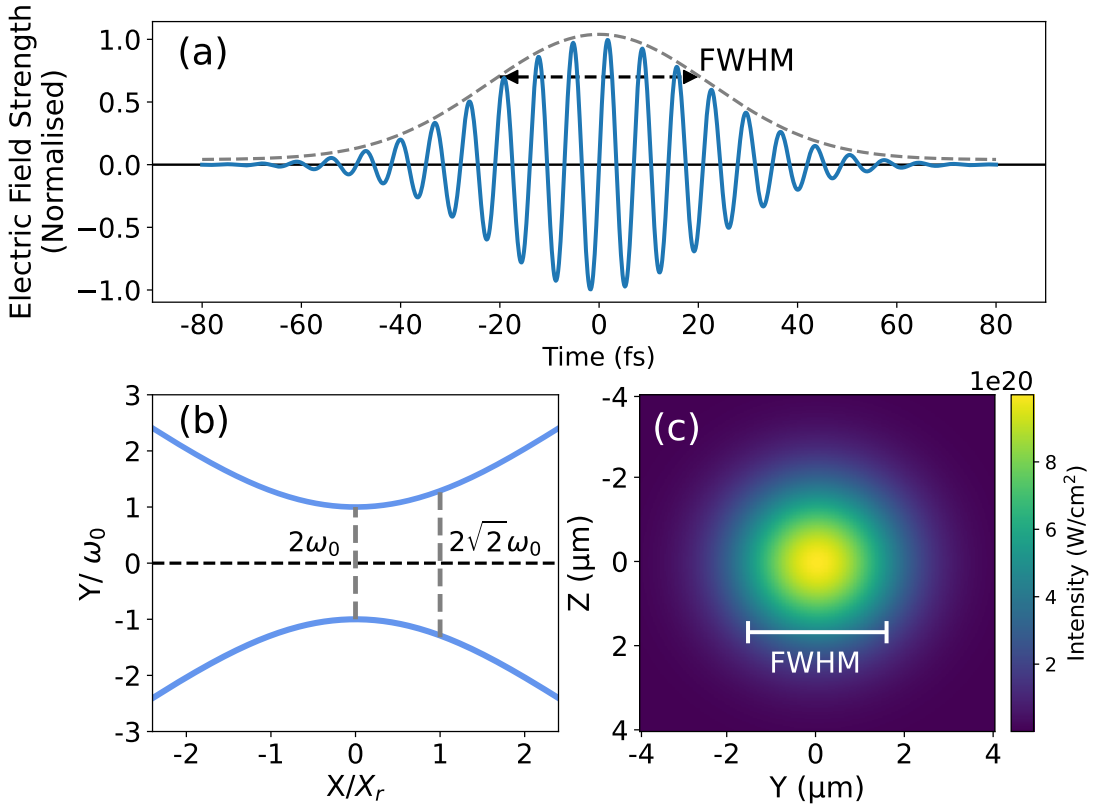


Figure 2.1: (a) A temporal intensity profile of a typical pulse used in a LPI, in which a sinusoidally varying electric field is seen, encapsulated within a Gaussian envelope. The FWHM of the Gaussian envelope is labelled, demonstrating that the pulse has a duration of $t_L = 40$ fs. (b) The focusing behaviour of a Gaussian beam, illustrating the minimum waist-size, $2\omega_0$, and the Rayleigh range, X_r . (c) An idealised transverse spatial intensity profile of a high intensity laser pulse at best-focus with the FWHM labelled.

Turning now to the term ‘high-intensity’, this typically refers to values ranging from 10^{18} – 10^{22} W/cm², with the upper limit often termed ‘ultra-high’ intensity. The values quoted when describing this aspect of a pulse will generally refer to the peak pulse intensity, achieved by bringing the initially collimated beam to focus with an optical component. In a typical experimental setup, this component would be an off-axis parabola (OAP) for which the focus position is offset from the center of the mirror, allowing the focused pulse to be directed away from the incoming beam. As the pulse is focused, it reaches a minimum diameter known as minimum waist, or ‘best-focus’, at which point the maximum intensity is attained. After the minimum waist diameter has been achieved, the beam expands symmetrically. This behaviour is illustrated in figure 2.1(b). In order to calculate the peak intensity of a pulse, the minimum waist or ‘spot size’ of the pulse’s spatial intensity profile must be measured. The peak intensity is of central importance when attempting to understand the physics of a given interaction due to the role that it plays in dynamics such as the ionisation of the target (section 2.2) and energy absorption by the resulting plasma (section 2.5). The spatial intensity profile at minimum waist is typically Gaussian in nature and is conventionally measured in one of two ways. The first is termed the waist-size, ω_0 (shown in figure 2.1(b)), given by

$$2\omega_0 = \frac{4\lambda_L f}{\pi D} = \frac{4\lambda_L F/\#}{\pi} \quad (2.1)$$

where D is the diameter of the beam incident on the focusing optic and f is its focal length, λ_L is the wavelength of the light being focused and $F/\#$ is the F-number of the optic. This is a metric commonly employed to characterise a focusing optic, defined as the ratio of the focal length of the optic to the diameter of the beam incident on it. Waist-size is defined as the distance over which the intensity of the pulse drops to $1/e^2$ of its maximum value. The second is simply the FWHM of the Gaussian profile, ϕ_L . These two metrics are related to one another through the expression $\phi_L = \omega_0\sqrt{2\ln 2}$. Figure 2.1(c) shows an example of an idealised transverse spatial profile of a beam at best focus, with the FWHM labelled. Throughout this thesis, it will be the FWHM measurement that is used when discussing focal spot size.

The final factor that must be taken into account when calculating the peak pulse intensity is the fraction of the pulse’s total energy, E_L , contained within the measured spot-size, η . For an idealised Gaussian spatial profile, this would be 0.5. However, a

typical realistic value for this would be between 0.3-0.4. This reduction arises from practical limitations in the focusing system, including wavefront aberrations, non-ideal beam quality and surface imperfections in the focusing optic.

These parameters can then be used with the following expression to calculate the peak intensity.

$$I_L = \frac{E_L \cdot \eta}{t_L \pi \frac{\phi_L^2}{2}} \quad (2.2)$$

For example, a typical pulse generated by a short pulse system such as Gemini could have the parameters $E_L = 10$ J, $\eta = 0.35$, $\phi_L = 4$ μm and $t_L = 40$ fs. These values would produce a peak pulse intensity of $I_L = 7 \times 10^{20}$ W cm^{-2} .

A final feature of a focusing beam that will be covered is that of the Rayleigh range, defined as the distance along the propagation axis of the pulse at which the intensity of the pulse is reduced by a factor of two, relative to the value at minimum waist size, given by

$$X_r = \frac{\pi \omega_0^2}{\lambda_L} \quad (2.3)$$

In waist size terms, this an increase by a factor of $\sqrt{2}$, illustrated in figure 2.1(b). In a practical context, knowing the Rayleigh range is beneficial as the target can be moved by a known distance to reduce the intensity, rather than using a different focusing optic, as altering this aspect of an experimental beam-line is often very time consuming.

Polarisation

The polarisation of light refers to the plane in which its electric field vector oscillates, with three main polarisation categories relevant to the research presented in this thesis: linear, circular and elliptical. Linear polarisation occurs when the electric field oscillates in a single plane, perpendicular to the pulse's propagation axis. For oblique laser-target interactions, two sub-categories of linear polarisations are relevant, termed 's' and 'p'. The first is the case of the electric field oscillating perpendicular to the plane of incidence, and the second is the case of the electric field oscillating parallel to that plane. Circular polarisation occurs when the magnitude of the electric field remains constant while its direction rotates in the plane transverse to propagation, such that the tip of the electric field vector traces a circle in time. Conceptually, this can be

considered as two linearly polarised pulses with equal magnitudes co-propagating 90° out of phase. Finally, elliptical polarisation is similar to the circular case, but with two co-propagating pulses of unequal magnitudes. This leads to the tracing of an elliptical path over time, with the pulse with the electric field strength of the greatest maximum amplitude referred to as the major axis.

2.1.2 Higher-order modes

Relevant to the results discussed in chapters 4, 5 and 4 is the generation of so-called higher-order modes (HOM). In laser optics, a mode is a transverse spatial intensity distribution that is preserved as the light propagates, aside from scaling and phase evolution. Mathematically, modes are defined as orthogonal solutions of the paraxial wave equation for the slowly varying field envelope U .

$$\nabla_{\perp}^2 U - 2ik \frac{\partial U}{\partial z} = 0 \quad (2.4)$$

The lowest-order solution is the fundamental mode, possessing a Gaussian intensity profile, discussed above and shown in figure 2.1(c). Higher-order modes correspond to other allowed solutions of the same equation and differ from the fundamental mode by possessing additional transverse structures, such as nodes or phase variations. In this thesis, the two forms of HOM of interest are Hermite-Gaussian (TEM) modes and Laguerre-Gaussian (LG) modes.

TEM modes are written with two indices as TEM_{mn} , where m and n are integers giving the number of transverse nodes in x and y . These indices appear in the Hermite polynomials and in the Gouy phase term that form part of the TEM solution to equation 2.4, given by

$$U_{mn}(x, y, z) \propto H_m\left(\frac{x}{w(z)}\right) H_n\left(\frac{y}{w(z)}\right) \exp\left(-\frac{x^2 + y^2}{w^2(z)}\right) \exp[i(m + n + 1) \zeta(z)] \quad (2.5)$$

where H_m and H_n are Hermite polynomials of order m and n , $w(z)$ is the beam radius at longitudinal position z , and $\zeta(z)$ is the Gouy phase.

LG modes are written as $\text{LG}_{p\ell}$ where p is the radial index and ℓ the azimuthal index. These indices appear in the Laguerre polynomial, the radial power term, the helical phase, and the Gouy phase, all present in the LG solution to equation 2.4, given

by

$$U_{p\ell}(r, \phi, z) \propto \left(\frac{r}{w(z)}\right)^{|\ell|} L_p^{|\ell|}\left(\frac{r^2}{w^2(z)}\right) \exp\left(-\frac{r^2}{w^2(z)}\right) \exp(i\ell\phi) \exp[i(2p + |\ell| + 1)\zeta(z)] \quad (2.6)$$

where $L_p^{|\ell|}$ is the associated Laguerre polynomial of radial order p and azimuthal order $|\ell|$, $w(z)$ is the beam radius at longitudinal position z , and $\zeta(z)$ is the Gouy phase.

2.2 Laser induced ionisation

Having discussed the characteristics of the laser pulses used in LPI's, we can now begin to cover the interactions themselves. Portions of the discussion in this section are adapted from [64]. The first aspect to cover is the generation of a plasma by the light preceding the main pulse from an initially solid density target. This light, composed of features known as the amplified spontaneous emission (ASE) pedestal and pre-pulse, typically occurs up to picoseconds prior to the main pulse and is discussed in more detail in section 3.2.2. While the exact details of what constitutes a plasma will be covered in a later section (2.3), at this point it is sufficient to know that it is a state of matter composed of atoms for which a number of electrons are no longer electrostatically bound to their parent nucleus. The process of freeing electrons from their parent nucleus, known as ionisation, occurs when an electron has gained sufficient energy to overcome what is known as their binding energy. For an initially solid-density target, this process requires the material to be heated by the laser pulse such that its constituent atoms are freed from the lattice structure, enabling ionisation to occur.

We begin by establishing a theoretical minimum pulse intensity at which an atom will be ionised by considering the simplest case; a single electron orbiting a single proton, i.e. the hydrogen atom. Understanding the ionisation dynamics at this basic level establishes a foundation on which the understanding of this process in more complex situations can be built. Calculating this threshold ionisation intensity for the hydrogen atom begins with its atomic radius, also known as the Bohr radius, a_B , given by

$$a_B = \frac{4\pi\epsilon_0\hbar^2}{m_e e^2} = 0.053 \text{ nm} \quad (2.7)$$

where m_e and e are the electron mass and charge respectively and $\hbar$ is the reduced

Planck's constant. From this radius, the electric field strength at the electron's orbital distance can be determined from the following equation;

$$\mathbf{E}_a = \frac{e}{4\pi\epsilon_0 a_B^2} \approx 5.1 \times 10^{11} \text{ V/m} \quad (2.8)$$

Lastly, from this electric field strength an equivalent intensity can be found from the following equation;

$$I_a = \frac{\epsilon_0 c E_a^2}{2} \approx 10^{16} \text{ W/cm}^2 \quad (2.9)$$

This value, often termed the atomic unit of intensity, states that any material in the presence of light with this intensity or greater will be ionised. However, experimental observations [65] and numerical modelling [66] show that ionisation can take place at intensities far lower than this, implying the existence of additional ionisation mechanisms, the details of which will now be covered.

2.2.1 Multiphoton ionisation

Multiphoton ionisation (MPI) [25, 67] is the process by which an electron absorbs a quantity of energy greater than its binding energy from multiple photons in a short period of time. This process is analogous to the photoelectric effect in which an electron is emitted from a material after absorbing energy from a single photon. However, given the wavelengths of the pulses generated by the systems used in the experiments detailed in this thesis (800 nm and 1.053 μm) the photon energies of 1.55 eV and 1.24 eV, respectively, are both below the binding energy of hydrogen (13.6 eV). In MPI, if the interaction between the electron and the laser field occurs over a timescale shorter than one laser cycle, the electron can absorb energy from multiple photons. This allows for the accumulation of enough energy, even from photons individually below the binding energy, for ionisation to occur. For example, given a pulse with a central wavelength of 800 nm, corresponding to a single photon energy of 1.55 eV, this would result in a hydrogen atom being ionised by 9 of these photons. However, it is often the case that an electron will absorb a greater number of photons than is required for ionisation, often termed above threshold ionisation [68]. This results in the electron possessing kinetic energy after ionisation that is given by

$$E_e = (n + s)\hbar\omega_L - E_b \quad (2.10)$$

where E_b is the binding energy of the atom, n the number of photons required to ionise the atom and s the number of surplus photons absorbed. MPI occurs at an intensity threshold of $\approx 10^{10}$ W/cm² and above. A schematic demonstrating this process is shown in figure 2.2(a).

2.2.2 Field ionisation by barrier suppression

Two additional relevant ionisation mechanisms exist, both of which can be understood by considering the potential barrier that an electron overcomes when ionisation occurs. This potential barrier is altered when an external electric field is applied, such as that present in high-intensity laser pulses. This scenario is analogous to an electron within a potential well, in which its ‘depth’ is altered by the presence of the laser field. The change in this potential well is given by

$$V(x) = \frac{-Ze^2}{x} - eE_{\text{ext}}x \quad (2.11)$$

where E_{ext} is the strength of the external electric field, x is the distance of the electron from the nucleus, and Z is the nuclear charge. As the laser pulse intensity increases, and consequently the strength of the external electric field, two distinct ionisation mechanisms arise. Both are forms of what is termed field ionisation by barrier suppression (FIBS), attributed to the external field’s suppression of the atomic field.

The first mechanism occurs when the applied external electric field is nearly equal in magnitude to the atomic electric field. In this case, the electron’s energy in the ground state, E_0 , is close to but slightly less than that required to escape the nucleus. Due to the quantum mechanical uncertainty in the electron’s energy, there is a finite probability that the electron can either pass over or ‘tunnel through’ the remaining potential barrier, leading to ionisation. This process is termed tunnelling ionisation, illustrated in figure 2.2(b).

The second mechanism arises when the external electric field exceeds the magnitude of the atomic field. In this situation, the electron’s energy in the ground state is sufficient to overcome the suppressed potential barrier, leading to ionisation. This is typically referred to as over-the-barrier ionisation (OBI). Tunnelling is not required for

the electron to escape from the nucleus as the potential barrier is fully suppressed. This mechanism is illustrated in figure 2.2(c).

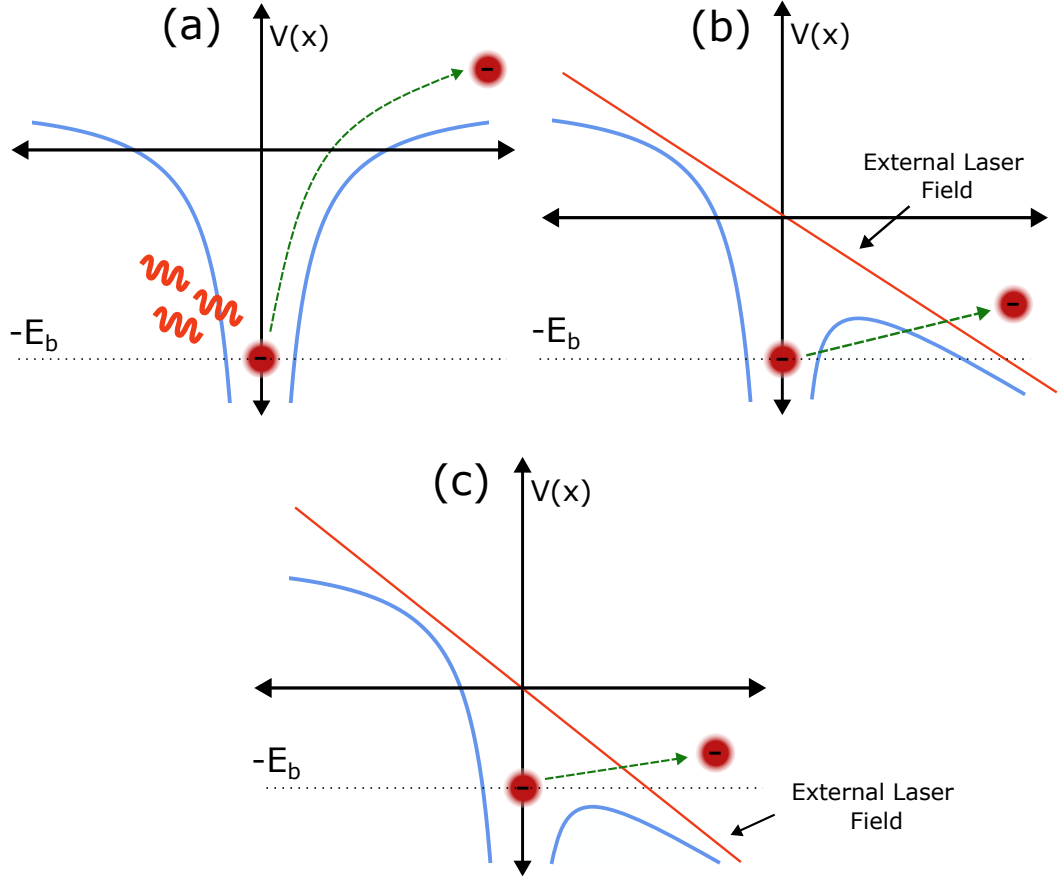


Figure 2.2: Illustrations of (a) Multi-photon, (b) tunnelling and (c) barrier suppression ionisation mechanisms, where E_b is the binding energy.

Given that for FIBS to occur, the atomic electric field must be suppressed, a corresponding laser electric field strength required for a given species of atom can be established through the following expression;

$$E_{FIBS} = \frac{4 \times 10^9 E_b^4 [eV]}{Z^2} \quad (2.12)$$

from which the pulse intensity at which these mechanisms occur can be found:

$$I_{FIBS} = \frac{\epsilon_0 c E_{FIBS}^2}{2} \quad (2.13)$$

Taking again the simplest example of the hydrogen atom, FIBS will begin to occur at intensities of $\approx 10^{14} \text{ W/cm}^2$.

A final useful metric concerning ionisation mechanisms is that of the so-called ‘Keldysh’ parameter [69]. This allows for the discernment of which of the forms of ionisation mechanisms, MPI or FIBS, is dominant at a given pulse intensity and frequency, given by

$$\gamma_K = \omega_L \sqrt{\frac{2E_b}{I_L}} \quad (2.14)$$

where ω_L and I_L are the laser frequency and intensity respectively. If γ_K is greater than 1, MPI is the dominant mechanism, if the inverse is true, FIBS is the dominant mechanism.

The ionisation mechanisms mentioned so far can be categorised as collisionless as ionisation occurs via the action of an external field, namely that of the laser pulse. However, there is an additional category of ionisation mechanisms known as collisional. Electrons that have already been freed from their parent nuclei can collide with bound electrons. If the kinetic energy of the freed electrons is sufficient, these collisions can cause ionisation through the imparting of this energy to the bound electrons. For this form of mechanism to become a relevant form of ionisation, there must be a large number of electrons already freed from their parent nuclei via collisionless mechanisms.

2.3 The plasma state

With an understanding of ionisation mechanisms established, the focus now turns to the properties of plasma to appreciate its role in interactions with high-intensity laser pulses. Plasma is often considered a distinct fourth state of matter, emerging through ionisation of atoms within a gas. However, a partially ionised gas does not automatically constitute a plasma. Broadly defined, such a gas must display two features to be considered a plasma, those being so-called ‘collective behaviour’ and quasi-neutrality. The first of these conditions is met when the behaviour of the ionised gas is determined not by short-range collisional forces, but rather by long-range electrostatic forces. A consequence of this is that a single particle can influence the behaviour of multiple others within the plasma. The second condition of quasi-neutrality is that the plasma is electrically neutral on a large enough scale, but not on a short scale. These two conditions can be understood through the examination of three commonly used criteria that will now be discussed.

The Debye length

The first criterion for a plasma is that its dimension, L must significantly exceed the so-called Debye length, λ_D , stated as $L \gg \lambda_D$. To understand the Debye length, consider the placement of a positive ‘test charge’ within the plasma. Being a collection of charged particles, the plasma’s electrons are quickly attracted to the positive test charge, while the positively charged ions are repelled. Due to their much higher charge-to-mass ratio, electrons respond to the test charge on a much shorter time-scale than the ions, forming a ‘cloud’ around it. This electron cloud effectively shields the surrounding plasma from the electrostatic potential of the test charge, limiting its influence to a finite region. The potential of the test charge in this setting no longer follows the typical inverse square law seen in free space but instead decays exponentially, as a function of the Debye length. The equation governing this modified potential is given as follows;

$$\phi(r) = \frac{q}{4\pi\epsilon_0 r} \exp\left(\frac{-r}{\lambda_D}\right) \quad (2.15)$$

where q is the charge of the test charge. The Debye length, λ_D is the distance over which the potential decays to a value of $1/e$ of its maximum value, itself given by

$$\lambda_D = \sqrt{\frac{\epsilon_0 k_B T_e}{e^2 n_e}} \quad (2.16)$$

where T_e and n_e are the electron temperature and number density, respectively, and k_B is Boltzmann’s constant. The Debye length can then be thought of as the distance characteristic of a given plasma beyond which charge neutrality is maintained.

Electron number density

The second criterion requires that the number of electrons within the cloud surrounding the test charge, N_D , within a Debye sphere with a radius equal to the Debye length, must be significantly greater than one. This requirement ensures that the plasma maintains quasi-neutrality and collective behaviour over larger scales. Mathematically, this condition can be expressed as follows:

$$N_D = \frac{4}{3} n_e \pi \lambda_D^3 \quad (2.17)$$

When this condition is satisfied, the plasma can effectively screen electrostatic potentials and exhibit the characteristic collective interactions that define plasma behaviour.

Plasma frequency

The third criterion states that the plasma frequency must be greater than the collisional frequency. The collisional frequency is simply the rate of collisions between the particles that comprise the plasma. The plasma frequency is a fundamental characteristic of a plasma, describing the oscillatory motion of the freed electrons within it. The cause of the plasma electron oscillations can be understood by considering the displacement of such electrons from the positively charged plasma ions. This charge separation establishes an electric field, attracting the ions and electrons toward each other. However, due to the equal charge magnitude carried by the two species of particles but large difference in mass, the electrons are accelerated at a much higher rate than the ions. This difference is so great that the ions can be considered to be effectively stationary on the timescales of the electron motion. The electrons are thus accelerated toward the stationary ions but overshoot them due to their inertia. This then establishes another charge separation and thus electric field, repeating the process of accelerating the electrons toward the ions. In this way, oscillatory motion of the electrons is established, with a characteristic frequency known as the plasma frequency, ω_p , given by

$$\omega_p = \sqrt{\frac{n_e e^2}{\epsilon_0 m_e}} \quad (2.18)$$

The relative frequencies establish which forces determine the dynamics of the ionised gas or plasma. If the collisional frequency is greater, it is hydrodynamic forces that are dominant. Whereas for the case of the plasma frequency being greater, it is electromagnetic forces that are dominant, leading to the collective behaviour described above.

2.3.1 Laser induced plasma expansion

An important aspect of any LPI is the plasma that is generated and expands from the front of a target during the initial stages of the interaction. While the cause of this expansion will be covered in section 3.2.2, the way in which this feature is quantified will be discussed due to the crucial role it plays in the dynamics of energy absorption

by the plasma from the laser pulse, discussed in the next section.

Briefly put, as the intensity of the pulse's rising edge increases throughout the initial stages of the interaction, a larger number of electrons are freed, leading to an exponentially decaying electron density profile being formed at the front of the target, as shown in figure 2.3.

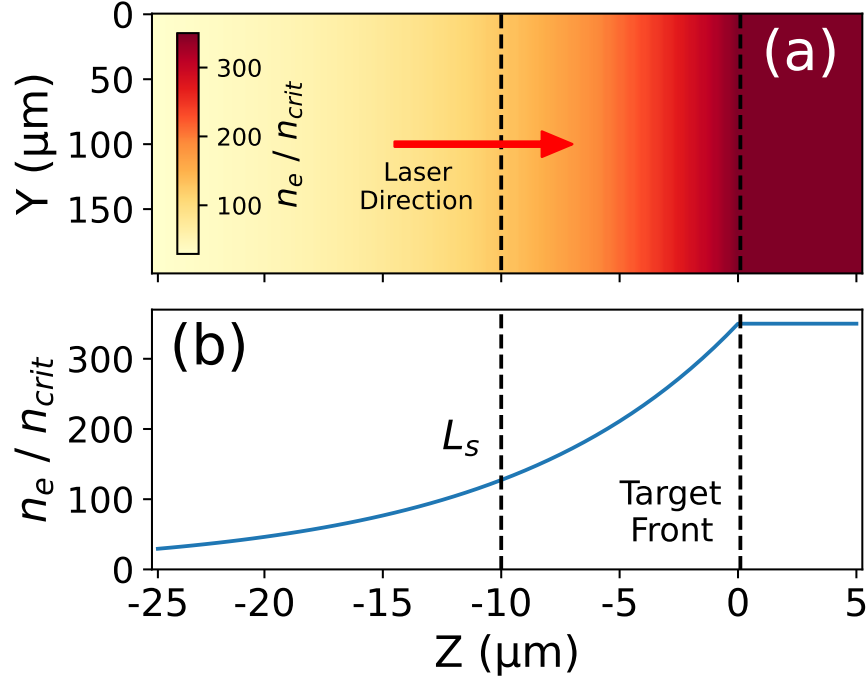


Figure 2.3: (a) A 2D map of the electron density typical of an LPI in which an exponential decay in density is seen with increasing distance from the target front. (b) A lineout taken from (a) at a position of $X = 100 \mu\text{m}$ showing the decaying electron density. Vertical lines showing the positions of the scale length and initial target front are present in both figures.

Note that the y -axis units used in figure (b) are normalised to n_{crit} , discussed below in section 2.4.3. This density profile is typically characterised by considering the description of the exponentially decaying electron density with the following expression;

$$n_e(z) = n_0 \exp\left(-\frac{|z|}{L_s}\right) \quad (2.19)$$

where n_0 is the initial target density, $n_e(z)$ is the electron density at a distance z from the initial target surface and L_s is what is known as its scale-length, that being the distance from the target at which the density has decayed to $1/e$ of its maximum value. It is the scale length that is used to quantify the expanded pre-plasma. Its value can be estimated using the following expression;

$$L_s = c_s t_L \quad (2.20)$$

where c_s is the local sound speed given by

$$c_s = \sqrt{\frac{Zk_B(T_e + T_i)}{m_i}} \quad (2.21)$$

where T_i is the ion temperature. However, the scale length is more commonly found through the use of hydrodynamic simulations (section 3.5.2), capable of simulating the intensity ranges and pulse durations that are responsible for the pre-plasma generation, or through optical probing of the front surface plasma during experiments [70][71]. Equation 2.19 is then fitted to the simulation results to determine the value of L_s .

2.4 Interaction between EM wave and plasma

With an understanding of both the formation and definition of a plasma covered, the next stage is to discuss the interaction between a high-intensity laser pulse and a plasma. Such an interaction consists of the complex interplay between EM fields of varying intensities and a plasma featuring both temperature and density gradients comprised of particles with different degrees of ionisation. To make possible the task of understanding the many dynamics at play in such a system, the approach is taken of first considering the simplified case of the interaction between a single electron and a plane wave. This case can then be built upon through the inclusion of more realistic features of a laser plasma interaction, thus ultimately giving insights into the physics being investigated in this thesis.

The physics of many important interaction dynamics can be understood through a consideration of only the motion of the electrons within a plasma in the presence of an EM field due to the quasi-stationary nature of the ions on the time-scales of the interactions being described. This is due to the differing charge to mass ratios of electrons and protons. As particles with equal magnitudes of charge but vastly different masses ($\frac{m_p}{m_e} = 1836$), it is evident that when acted upon by electric and magnetic fields they will experience equal magnitudes of forces but much different rates of acceleration. It is for this reason that electron dynamics will be the main focus in this section.

2.4.1 Electron motion in an infinite plane wave

The simplest case of the motion of an electron due to the influence of an EM wave is that of an infinite plane wave, that being a wave with a constant maximum amplitude and infinite extension in space. In the presence of such a field, the force exerted on an electron due to both the electric and magnetic fields is given by the Lorentz force equation as follows;

$$\mathbf{F}_L = \frac{d\mathbf{p}}{dt} = -e(\mathbf{E} + \mathbf{v} \times \mathbf{B}) \quad (2.22)$$

where $\mathbf{E}$ and $\mathbf{B}$ are the electric and magnetic field, respectively, $\mathbf{p} = \gamma m_e \mathbf{v}$ is the momentum and $\gamma = \sqrt{1 + v^2/c^2}$ is the relativistic Lorentz factor. An important element in determining the motion of the electron is the $1/c$ factor difference between the magnetic and electric field strength magnitudes. Given that the component of the force exerted on the electron by the magnetic field is found through its cross-product with the electron velocity $\mathbf{v}$, it follows that for the magnitude of this force to be non-negligible in comparison to that exerted by the electric field, the electron must be travelling at near relativistic velocities. This then provides two regimes under which the motion of an electron in an infinite plane wave can be considered (i) non-relativistic velocities ($v \ll c$) and (ii) relativistic velocities ($v \approx c$).

In regime (i), the $\mathbf{B}$ component is negligible in the Lorentz equation, and so equation 2.22 reduces to

$$\frac{d\mathbf{p}}{dt} = -e\mathbf{E} \quad (2.23)$$

Then considering the form of the $\mathbf{E}$ field oscillation in a plane wave, given by

$$\mathbf{E} = E_0 \sin(\omega_L t - ky) \hat{\mathbf{x}} \quad (2.24)$$

where $k = \omega_L/c$ is the wave number. Equations 2.23 and 2.24 can be combined to describe the acceleration of an electron under the effect of such a field

$$\frac{\partial v}{\partial t} = \frac{-eE_0}{m_e} \sin(\omega_L t - ky) \quad (2.25)$$

which when integrated with respect to time gives the velocity of the electron as it oscillates in the electric field, also known as the quiver velocity, given by

$$v_{\text{osc}} = \frac{eE_0}{m_e\omega_L} \cos(\omega_L t - ky) \quad (2.26)$$

The maximum quiver velocity is achieved when the electron is at its midway point of oscillation, given by

$$v_{\text{max}} = \frac{eE_0}{\omega_L m_e} \quad (2.27)$$

With sufficient magnitude, the **E** field can accelerate the electron to relativistic velocities such that the quiver velocity approaches c , at which point the motion in regime (ii) can be considered. With the effect of the **B** field component no longer negligible, it must be included in the Lorentz force equation. The electric and magnetic field are related by the following expression

$$|B_0| = \frac{1}{c}|E_0| \quad (2.28)$$

which can be substituted into equation 2.22 to describe the force exerted on an electron by both fields. In a similar fashion to the equations presented above, the integral with respect to time can be taken to give an expression for the velocity of an electron due to both fields

$$v_{\text{tot}} = \frac{e^2 E_0^2}{4m_e^2 \omega_L} \left[\cos(\omega_L t - ky) + \frac{1}{2} \sin(2\omega_L t - 2ky) \right] \quad (2.29)$$

The motion in the transverse direction is unchanged from that occurring under regime (i), as the $\cos(\omega_L t - ky)$ component is equal to that seen in equation 2.26. However, the addition of the $\sin(2\omega_L t - 2ky)$ component describes electrons being accelerated in the wave propagation direction with a frequency twice that of the laser. The effect of both the **E** and **B** fields combine to cause a figure-of-eight motion of the electron as it propagates

The two forms of electron motion, while in some sense idealised due to the infinite plane wave that causes the motion, can have significant effects on the interaction dynamics, such as absorption of laser energy into the plasma, covered later in section 2.5. As such, the step can be taken to normalise the quiver velocity of the electrons to c , with the parameter known as the normalised light amplitude, given by

$$a_0 = \frac{v_{\text{osc}}}{c} = \frac{eE_0}{m_e \omega_L c} \quad (2.30)$$

If $a_0 \ll 1$, electrons in the presence of the electric field with strength E_0 are not relativistic, and are therefore considered in the regime of motion (i). Whereas, if $a_0 \geq 1$ they reach relativistic velocities and thus their motion is understood through regime (ii). In more practical terms, the fact that this is related to the electric field strength allows for relativistic motion of electrons to be related to the intensity of a laser pulse, providing a threshold intensity, dependent on pulse wavelength, above which electrons will begin travelling at relativistic velocities, given by

$$I_L = \frac{a_0^2}{\lambda_L^2 [\mu\text{m}]^2} \times 1.37 \times 10^{18} \text{ W/cm}^2 \quad (2.31)$$

For the two laser systems used in the experimental campaigns reported on in chapters 6 and 7 (Gemini and Vulcan, with $\lambda_L = 800 \text{ nm}$ and $1.053 \mu\text{m}$ respectively) the intensity thresholds for achieving relativistic electron velocities are $2.1 \times 10^{18} \text{ W/cm}^2$ and $1.2 \times 10^{18} \text{ W/cm}^2$, respectively.

2.4.2 Electron motion in a non-homogeneous EM field

While analysing the case of an infinite plane EM wave provides valuable insight into fundamental aspects of electron motion, it does not include critical features of real laser pulses used in typical interactions, namely the spatial and temporal variations in field strength. As discussed in section 2.1.1, these variations are typically Gaussian in nature. To understand the impact of non-homogeneous EM fields, consider the forces acting on an electron initially at rest during a single laser cycle. The oscillating electric field will exert a force that accelerates the electron from rest, displacing it from its initial position. As the field reverses direction, a ‘restoring force’ is applied, accelerating the electron back toward its initial position. In the case of a uniform field, these forces are equal in magnitude, resulting in the electron returning to its original position, shown in figure 2.4(a). However, when the field strength is non-uniform, these forces become imbalanced, preventing the electron from returning to its initial location. Instead, the electron’s final position will shift toward regions of lower field strength.

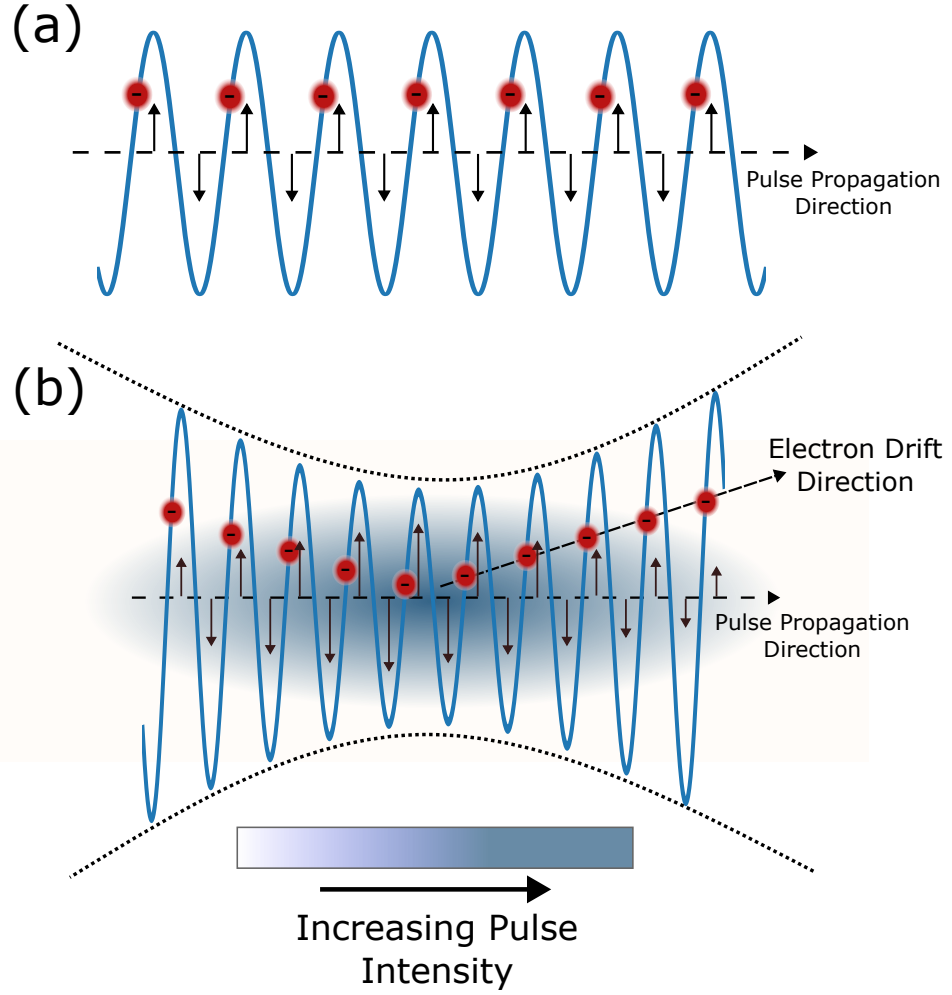


Figure 2.4: Illustration showing the path of an electron over time in a (a) uniform EM field and a (b) non-uniform field such as that produced by a laser pulse being brought to focus. The arrows represent the magnitude of the restoring force exerted on the electron.

Over multiple laser cycles in a non-uniform field, this field strength imbalance results in the net application of a force causing a ‘drift’ of the electron toward regions of lower intensity shown in figure 2.4(b). Considering the Gaussian spatial profile of a typical laser pulse used in an LPI, this phenomenon leads to electrons being pushed away from the central axis towards the ‘wings’ of the pulse. This force, termed the ponderomotive force [64], can be described by considering the gradient of the electric field strength. As with the case of the infinite plane wave, the electron quiver velocity determines the nature of the electron’s motion, leading to two regimes. In the non-relativistic case ($v_{osc} \ll c$), the ponderomotive force is given by

$$\mathbf{F}_p = -\frac{e^2}{4m_e\omega_L^2}\nabla\mathbf{E}^2 \quad (2.32)$$

In the second regime for which $v_{osc} \approx c$, and for linearly polarised light, this force can be described in terms of the normalised light amplitude by the following expression:

$$\mathbf{F}_p = -m_e c^2 \nabla \sqrt{1 + a_0^2/2} \quad (2.33)$$

2.4.3 Non-homogeneous EM field interacting with multiple electrons

As discussed in previous sections, the leading edge of a high-intensity laser pulse interacting with a target results in the freeing of a significant number of electrons from the atoms that comprise the target. This process, facilitated by the ionisation mechanisms previously described, produces a collection of free electrons that constitutes a plasma, provided certain conditions are satisfied (2.3).

This section is concerned with the interaction between the remainder of the pulse and this laser-generated plasma, with a focus on the effect the plasma has on the propagation of the pulse. We begin by deriving the wave equation governing the behaviour of an electromagnetic wave as it propagates through a plasma, given by the following expression:

$$\nabla^2 \mathbf{E} = \frac{1}{c^2} \frac{\partial^2 \mathbf{E}}{\partial t^2} + \epsilon_0 \frac{\partial \mathbf{J}}{\partial t} \quad (2.34)$$

where $\mathbf{J} = -en_e \frac{dr}{dt}$ is the current density, present due to the plasma consisting of charged particles. From this, a dispersion relation can be derived in which the plasma frequency and the laser frequency are related to one another.

From this, the dispersion relation for an EM wave propagating through plasma can be derived, given by

$$\omega_L^2 = c^2 k^2 + \omega_p^2 \quad (2.35)$$

where ω_p is the plasma frequency. The refractive index of a plasma can be found from the relation $\eta = \frac{ck}{\omega_L}$, where k is the wave number. This then gives the following expression for the refractive index

$$\eta = \sqrt{1 - \frac{\omega_p^2}{\omega_L^2}} \quad (2.36)$$

This expression demonstrates that if ω_p is greater than ω_L , the refractive index is imaginary, establishing the condition of $\omega_L > \omega_p$ for the remainder of the pulse to be able to propagate through the plasma. What is occurring physically is that as the pulse begins to interact with the plasma, the oscillating electrons respond to the oscillations of the electric field. In the case of the plasma frequency being greater than the pulse frequency, the electrons respond on a timescale shorter than the field oscillations, effectively dampening them and so preventing propagation of the pulse. It is informative to examine both the group and phase velocities of an EM wave as it propagates through a plasma, which can be derived from equation 2.35. Beginning with the phase velocity, $v_p = \frac{\omega_L}{k}$, this is given by

$$v_p = \sqrt{c^2 + \frac{\omega_p^2}{k^2}} \quad (2.37)$$

This expression demonstrates that for any EM wave that is propagating through a plasma, the phase velocity must be greater than c , appearing to be in violation of the laws of relativity. However, this can be resolved by considering that what is being described by this expression is the velocity of the sinusoidal oscillations of the wave, not the variations in its amplitude within which information can be contained, thus no laws of relativity are being violated.

The group velocity, $v_g = \frac{\partial \omega_L}{\partial k}$, is given by

$$v_g = \frac{c^2}{\sqrt{c^2 + \frac{\omega_p^2}{k^2}}} = \frac{c}{\sqrt{1 + \frac{\omega_p^2}{c^2 k^2}}} \quad (2.38)$$

This demonstrates that for any EM wave propagating through a plasma, its group velocity must be less than c . The group velocity describes the velocity of the varying amplitude of the wave, which as mentioned above, contains information.

Having now established that the relative frequencies of the pulse and plasma plays a pivotal role in determining whether the plasma will be opaque to the pulse or not, the factors influencing these metrics can be considered. Equation 2.18 demonstrates the dependence of the plasma frequency on the electron number density of the plasma. By setting the plasma frequency equal to the laser frequency, this expression can be

rearranged to give the electron number density above which the plasma is opaque to the laser light. This is known as the critical density, n_{crit} , and is given by

$$n_{crit} = \frac{m_e \epsilon_0 \omega_L^2}{e^2} \quad (2.39)$$

This expression then allows for the defining of two distinct regimes of plasma density, termed overdense and underdense. The first being when $n_e > n_{crit}$ and the plasma is opaque to the laser light, and the latter being when $n_e < n_{crit}$ and the plasma is transparent to the laser light.

The effect of these two regimes can be seen by considering the pre-plasma formed at the target front prior to the arrival of the main pulse (2.3.1). The electron number density decays exponentially with distance from the target, so as the pulse propagates through this pre-plasma, it will do so until reaching an electron number density for which $n_e = n_{crit}$. At this surface within the pre-plasma, known as the critical surface, it will be reflected by the overdense plasma.

However, the response of the plasma electrons to the EM fields that prevents the pulses propagation is not instantaneous. This allows for the pulse to propagate into the opaque plasma for a short period of time, decaying exponentially with distance. This can be characterised by what is termed the skin depth, defined as the distance over which the pulse amplitude decays to a magnitude of $1/e$ of its maximum value prior to interaction with the plasma. This metric, δ , is determined by the relative frequencies of the plasma and light by the following expression;

$$\delta = \frac{c}{\sqrt{\omega_p^2 - \omega_L^2}} \quad (2.40)$$

2.4.4 Relativistic self-induced transparency

The final optical effect that will be discussed is that of relativistic self-induced transparency (RSIT) [72], a phenomenon particularly relevant to the results presented in chapter 7 of this thesis in which the formation of so-called ‘relativistic plasma apertures’ [57] is studied. This effect allows for the propagation of a laser pulse into a plasma with electron densities that are higher than would classically be allowed. This is attributed to the relativistic mass increase of the electrons if they are accelerated to velocities approaching c . Equation 2.18 can be modified with the Lorentz factor to

account for this increase in electron mass, leading to a decrease in plasma frequency, as described by

$$\omega'_p = \sqrt{\frac{n_e e^2}{\epsilon_0 \gamma m_e}} \quad (2.41)$$

This then leads to an increase in the critical density described by

$$n'_{\text{crit}} = \gamma n_{\text{crit}} = \frac{\gamma m_e \epsilon_0 \omega_L^2}{e^2} \quad (2.42)$$

This can impact interaction dynamics in multiple ways [58], such as moving the critical surface further into the target [73], altering the laser energy absorption mechanism [74], to entire regions of a target becoming completely transparent to laser light, allowing for larger percentages of the pulse being transmitted than would otherwise occur [75].

2.5 Laser-to-electron energy absorption mechanisms

The final aspect of an interaction that will be discussed is that of the transfer of energy from the laser pulse to the plasma electrons via so-called absorption mechanisms. The generation of light from LPI's is central to the work reported on in this thesis, and the transfer of energy from the drive laser pulse to this generated light is mediated by the plasma electrons, hence the importance of these generation mechanisms. These can be broadly divided into two categories termed collisional and collisionless, with the form of mechanism that is primarily responsible for energy transfer depending on the pulse intensity. Collisional mechanisms are relevant up to intensities of $\approx 10^{15}$ W/cm², while above this level, it is the collisionless mechanisms that are dominant.

2.5.1 Collisional absorption

As mentioned above, this form of absorption mechanism is dominant for intensities up to $\approx 10^{15}$ W/cm², levels which are primarily present in the ASE pedestal and the rising edge of the main pulse (discussed in section 3.2.2) and the lower intensity edges of the pulses transverse spatial profile. While this mechanism is not primarily responsible for phenomena such as RSIT or electron acceleration—two factors likely to influence the generation of light during an interaction—it is responsible for the initial heating of the

bulk of the target prior to the arrival of the higher intensity portion of the pulse and thus warrants discussion.

The most relevant collisional absorption mechanism is that of inverse bremsstrahlung [76]. As mentioned above, an electron will experience an energy gain as it is accelerated in the first half of the laser-cycle. However, if a collision between such an electron and a plasma ion occurs, a portion of the electron's energy will be transferred to the ion which cannot then be recovered by the fields in the second half of the laser-cycle. A key factor is that of the mean-free-path of the electrons, which in this intensity regime is less than the target dimensions, allowing for collision between electrons and ions to occur.

2.5.2 Collisionless absorption

Moving now onto the dominant form of absorption mechanism, that being collisionless. As the mean-free-path of the electrons increases beyond that of the plasma dimensions, the primary way in which an electron transfers its energy to the plasma is by its propagation past the plasma critical surface. This is facilitated by the increased magnitude of the laser electric field strength at intensities $> 10^{15}$ W/cm², relative to the lower EM strengths at which collisional absorption dominates, and as such the oscillations of the electrons are large enough to move past the critical surface as required. The mechanisms most dominant in the interactions being studied in this thesis are as follows.

Resonance absorption

The first absorption mechanism to be covered is termed resonance absorption and occurs when a p-polarised pulse is obliquely incident upon a plasma density gradient at the target front for which the scale length is greater in magnitude than the laser wavelength, $L_S > \lambda_L$. As the pulse propagates into denser plasma, it will reach a point at which it will be reflected, resulting in the direction of the electric field component of the pulse being normal to the density gradient. The field extends beyond the turning point to the critical density leading to the excitation of plasma density waves that transfer energy to electrons in this volume of the overdense plasma, shown in the schematic in figure 2.5(a). This process occurs most efficiently at the point of the pulse's reflection, when the plasma and laser frequencies are equal (resonant), hence the mechanism is termed

resonance absorption.

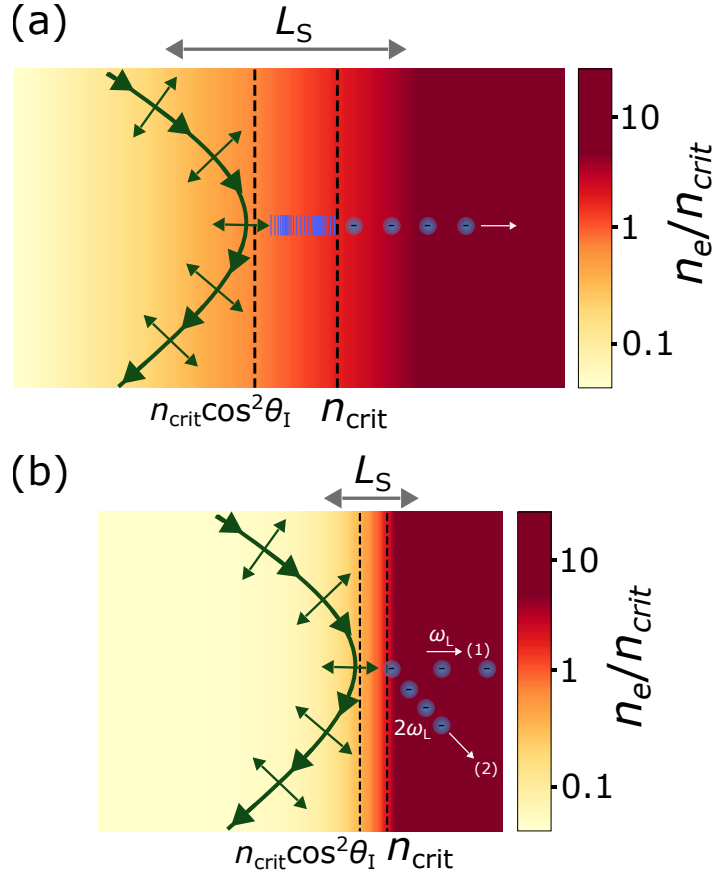


Figure 2.5: (a) An illustration of the resonance absorption mechanism in which a p-polarised pulse is obliquely incident on the pre-plasma at the target front for which $L_S > \lambda_L$. (b) An illustration of the vacuum heating mechanism (1) in which a p-polarised pulse is obliquely incident on a pre-plasma at the target front for which $(L_S < \lambda_L)$. Also illustrated is the relativistic $\mathbf{j} \times \mathbf{B}$ heating mechanism (2).

An important consequence of oblique incidence is a modification of the plasma density at which the pulse is reflected. As the laser propagates into a plasma with a density gradient, the refractive index varies continuously, causing the light to refract and bend progressively toward the target surface. The result is that the density at which the pulse is reflected is lower than if the angle of incidence were 0, being expressed by the expression $n_{\text{crit}} \cos^2 \theta_I$.

At the point of reflection, when the polarisation plane of the pulse is normal to the density gradient, electrons are oscillated past the critical surface. Due to the exponential decay of the $\mathbf{E}$ field strength into the plasma, there is a weaker restoring force applied to these electrons, and as such they are not drawn back into the underdense plasma, therefore experiencing a net gain in energy. This leads to a build up of electrons

and thus a charge separation, exciting electron density waves in this region of the plasma. These waves will propagate until a phenomenon known as wave-breaking occurs [77]. As the amplitude of an induced plasma wave increases over multiple laser cycles, the resulting increase in charge separation leads to higher velocities of the oscillating electrons that constitute the plasma wave. Wave-breaking occurs when the electron velocities begin to exceed that of the phase velocity of the plasma wave such that its coherent structure cannot be maintained. Energy present in the plasma wave is then transferred to the electrons in this region of the target. This may occur over multiple wave cycles with a lower intensity pulse, or over a single cycle if the pulse is sufficiently intense.

Vacuum heating

Vacuum heating [78] occurs under similar conditions to resonance heating, when a p-polarised pulse is obliquely incident onto the target front but with a key difference of a much shorter scale length, $L_S < \lambda_L$. As the pulse is reflected, the electric field draws electrons from beyond the critical surface accelerating them into the vacuum, hence the name of this mechanism. As the force initially responsible for their acceleration was provided by the exponentially decaying electric field present past the critical surface, the restoring force applied to the electrons in the vacuum is much greater. Thus they are accelerated into the plasma, past the critical surface beyond the influence of the exponentially decaying restoring force, leading to a net gain in their energy. This process is illustrated in figure 2.5(b). The much sharper plasma boundary in this case means that plasma waves cannot be supported, preventing resonance absorption from occurring. This results in electrons being injected into the plasma at roughly the quiver velocity, once per laser cycle.

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

The final absorption mechanism relevant to the work presented in this thesis is termed relativistic $\mathbf{j} \times \mathbf{B}$ heating [79]. This mechanism becomes dominant at relativistic pulse intensities. As discussed in section 2.4.1, as a_0 approaches 1, the magnitude of the $\mathbf{B}$ field is no longer negligible, and so begins to influence the electron's motion as described by the $\mathbf{v} \times \mathbf{B}$ component of equation 2.22. This force is described by

$$F = -\frac{m_e}{4} \frac{\partial v_{os}^2(x)}{\partial x} (1 - \cos 2\omega_L t) \quad (2.43)$$

The presence of the $2\omega_L$ term indicates that electrons are accelerated past the critical surface twice per laser cycle, and therefore twice the rate of the other absorption mechanisms discussed. In addition, this force is directed along the laser propagation axis, notably different from the other absorption mechanisms discussed so far for which this has been in the direction of the electric field. The effect of this force is that electrons are accelerated past the critical surface such that they are no longer subject to the influence of the laser fields, thus experiencing a net gain in energy. This mechanism is illustrated in figure 2.5(b) in which the differences in electron acceleration direction and frequency relative to the vacuum heating mechanism are shown.

This mechanism is also strongly dependent on pulse polarisation. It occurs for linear polarisation, both ‘s’ and ‘p’, but does not occur for circularly polarised light as the magnetic field component averages to 0 in equation 2.29.

2.6 Summary

To summarise, this chapter has provided an overview of the key physics concepts necessary to understand the results presented in Chapters 1, 2, and 3. It begins with an overview of the high-intensity pulses commonly used in LPI’s, followed by a discussion of how these pulses ionise the target surface to create a plasma. The dominant ionisation mechanism is shown to be dependent on the intensity and frequency of the laser pulse. From there, the chapter explores the definition of a plasma, outlining the characteristics of quasi-neutrality and collective behaviour, both of which must be present for matter to be classified as a plasma. The discussion then moves to the formation of the expanded plasma at the front of the target and how this is characterised by the scale length. Next, the interaction between an EM wave and plasma is examined. This begins with the simplest case of an electron interacting with an infinite plane wave, followed by consideration of the more complex dynamics present in realistic interactions. Finally, the chapter discusses the various mechanisms by which energy from the laser pulse is absorbed by the plasma. These mechanisms are shown to depend strongly on the pulse’s polarisation, angle of incidence, and the scale length of the pre-plasma.

Chapter 3

Methodology

This chapter provides an overview of the key experimental and numerical simulation techniques that were employed to investigate novel aspects of laser-plasma interaction physics that are reported on in this thesis. This will begin with a description of how a typical experimental campaign in this field is conducted, after which an overview of the technology that allows for the generation and use of high intensity ($> 1 \times 10^{18}$ W/cm²) pulses at the desired repetition rate will be discussed. Next, details of the two high-power laser facilities used to conduct the experiments reported on in this thesis, Gemini and Vulcan, will be covered, followed by a discussion of the main diagnostic techniques that were used during these experiments. Finally, the forms of numerical simulation that were employed to aid in the interpretation of the measurements made during the experimental campaigns to understand the underlying physics will be explored, namely particle-in-cell (PIC) and hydrodynamic.

3.1 Experimental overview

The technology required to generate laser pulses with the necessary parameters for investigating the physics in this field is highly specialised and exists at both universities and national-scale laboratories. Access to such facilities by external research groups is a limited resource and can only be gained through proposals submitted by said groups in which the aims of the experiment and how they will be achieved are detailed.

Once facility access time, often referred to as beam-time, has been granted, typically between 3 and 6 weeks in duration, a period of several months is spent planning the experiment. This is a multi-stage iterative process that involves both staff and students

from the research group that will be conducting the experiment and staff based at the research facility itself. The plans made during this period are vital in determining the success of the experimental campaign, ensuring that the aims are realistic and strategies for overcoming potential obstacles are developed. The experiment is typically carried out by between 6 and 8 person, including students, post-docs and researchers, each assigned various responsibilities and tasks. The overall direction of the experiment is determined by a senior member of the research group known as the ‘principal investigator’ or PI. The activities that occur within the target area itself are typically managed by senior members of the research group. The experiment begins with a setup period, typically one or two weeks, during which the interaction point between target and laser pulse is defined within the interaction chamber. Once this point has been defined, the mechanical and optical infrastructure that will facilitate the experiment, such as target drive and focal spot cameras, can be put in place and the diagnostics can be aligned to the interaction point. At this point the aims of the experiment can be reassessed to ensure they are realistic, as the conditions present may require minor changes to the objectives. Once the setup has been completed, the diagnostics are calibrated and optimised to ensure the validity of the measurements that will be made. This is typically performed by firing initial shots with the laser system which has the added benefit of a ‘shake-down’ of the system itself to ensure that it is operating as expected. Once these tasks have been completed, the experiment will shift into a data collection mode during which the remaining beam-time is devoted to completing the aims of the experiment by collecting high-quality data that will hopefully uncover new physics.

3.2 High power laser technology

When conducting research in the field of laser-plasma interactions, the characteristics of the laser pulse being used are of the utmost importance. As these characteristics are determined by the particular laser facility being used, an understanding of the work presented in this thesis will be aided by having a working knowledge of how high-power laser systems function. While their engineering and design are beyond the scope of this thesis, this section aims to cover the principles of some of the technologies that are most important to the generation and use of high-intensity laser pulses.

3.2.1 Pulse generation and amplification

This section will give a brief overview of the way in which the high-power laser systems used to obtain the results presented in this thesis generate relativistically intense laser pulses. While the exact way in which this is achieved varies between individual systems, the technique common to both that were employed is that of chirped pulse amplification (CPA) [10], the advent of which allowed for the generation of peak laser pulse intensities exceeding 10^{18} W/cm².

The CPA process begins with the generation of a short-duration (tens of fs) low energy (nJ) ‘seed’ pulse from an oscillator. The duration of the pulse is then increased, or stretched, by passing it through a pair of optical components known as diffraction gratings. These components rely on a feature of the seed pulse known as its bandwidth, which refers to the range of wavelengths it contains. As the pulse passes through the diffraction gratings, different path lengths are created for the range of wavelengths present in the pulse, inducing what is known as a group velocity dispersion. This increase in duration is referred to as chirping, hence the technique is named chirped pulse amplification.

The next stage of the CPA process is the amplification of the now stretched pulse. Various amplification techniques, such as regenerative and multi-pass, are utilised depending on the specific laser system in use. The fundamental principle of these techniques involves a material known as a gain medium into which energy is optically pumped using a light emitting component, for example a flash lamp or light emitting diode. The transfer of this energy results in an excitation of the electrons within the gain medium from their ground states into higher energy levels. When the majority of electrons in the gain medium occupy this higher energy state, population inversion has been said to have been achieved. The pulse that is undergoing amplification is then propagated through the population inverted gain medium. This causes the excited electrons to transition back down to their lowest energy level in a process known as stimulated emission, releasing photons as they do so. The energy of the photons, and hence their wavelength, is determined by the difference in energy between the excited and ground states of the material comprising the gain medium. This material is chosen such that the wavelength of the emitted photon matches that of the pulse being amplified. As a result, the photons generated via stimulated emission contribute to the

existing photons in the pulse, thereby amplifying it.

It is the amplification stage that highlights the crucial role that the CPA technique plays in the generation of high-intensity pulses. Had the seed pulse been amplified directly, its intensity would rapidly reach a level ($\sim 10^{10}$ W/cm²) that would damage the solid-state optics that comprise the laser-chain, preventing further amplification. In addition, damage to optical components can occur from nominally lower intensity pulses due to non-linear effects. Examples of such effects are those that stem from the optical Kerr effect. In a medium in which this effect manifests, the refractive index is dependent on the intensity of the light as it propagates through the material, described by the following expression.

$$\eta = \eta_0 + \eta_2(I_L) \quad (3.1)$$

where $\eta_2(I_L)$ represents the intensity dependent component.

A consequential non-linear effect can be understood by considering the transverse spatial profile of the pulse. As discussed in the previous chapter, the pulses used in this field typically have a Gaussian transverse spatial distribution. The variation in intensities across this profile results in a corresponding variation in refractive index, with the lower intensity edges of the pulse propagating through the medium faster than the higher-intensity center. This results in a focusing of the pulse to higher intensities in a phenomenon termed self-focusing. Consequently, the damage threshold of solid state optics can be exceeded for pulses that are nominally below this limit. The temporal variations in the pulses intensity give rise to an additional relevant non-linear effect. As the pulses intensity changes over time, a corresponding time-dependence in the refractive index that it experiences occurs. This phenomenon, known as self-phase modulation, can lead to a broadening of the pulse bandwidth and an increase in its duration. Non-linear effects can be quantified by the ‘B-integral’, describing the accumulated phase after the pulse has propagated through an optical medium a distance of L ;

$$B = \frac{2\pi}{\lambda_L} \int_0^L \eta_2 I_L(z) dz \quad (3.2)$$

with values of $B \approx 1$ -2 rad recommended for avoiding such deleterious effects [80].

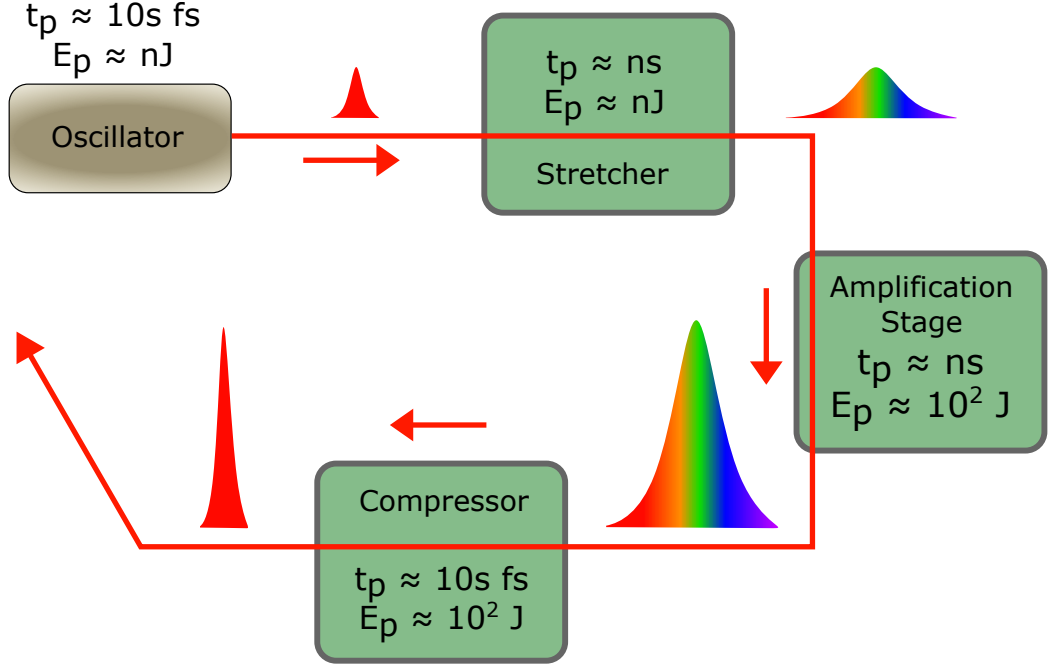


Figure 3.1: Schematic showing a simplified form of the CPA technique, displaying the typical pulse durations, t_p and pulse energies E_p that could be expected after each stage.

An approach that could be used to ensure that the pulse does not exceed the damage threshold is to increase the beam diameter. However, the cost of the numerous large-diameter optics that would be required makes this approach prohibitively expensive. Instead, the approach of chirping the pulse is used, reducing the pulse intensity such that even after the increase in energy that the amplification stage provides, the damage threshold is not exceeded. Once amplification is complete, the pulse can contain energies between tens and hundreds of Joules. The final stage of CPA is reversing the chirp that was induced in a process known as pulse compression. This reduces the duration of the pulse to femtoseconds or picoseconds, with the minimum achievable duration being dependant on the pulse bandwidth, as described by;

$$\Delta\nu\Delta t \geq K \quad (3.3)$$

where $\Delta\nu$ is the pulse bandwidth, Δt is the pulse duration and K is a pulse shape dependent constant, which for a Gaussian pulse shape has a value of 0.44.

The resulting high-energy, short-duration pulse can then be focused to relativistic intensities. A schematic detailing the various stages of the CPA process is shown in figure 3.1.

3.2.2 Temporal intensity contrast

When the pulse temporal intensity profile was discussed in the previous chapter (section 2.1.1), a somewhat idealised version was presented. In reality, the main Gaussian profile of the pulse is preceded by lower intensity light that can exist on the timescales of picoseconds to nanoseconds. While the intensity of this light is orders of magnitude lower than that of the main pulse, when focused by the low F-number optics used to achieve peak pulse intensities on the order of 10^{21} W/cm², this preceding light can reach intensities that exceed the ionisation threshold ($\sim 10^{10}$ W/cm²). The resulting plasma generation and target expansion can significantly alter interaction dynamics. For example, the dominant absorption mechanism can be altered through changes to the pre-plasma scale length (section 2.5), or expansion over the timescales of nanosecond can result in the entire target becoming underdense or being destroyed prior to the arrival of the main pulse [81].

A metric that is commonly employed to quantify this preceding light is the temporal intensity contrast of the pulse, more commonly referred to as the pulse contrast. This is defined as the ratio of the peak intensity of the main pulse to that of the light preceding it. Contrast measurements are quoted with the time relative to the pulse peak at which the background laser light is sampled. This is due to the preceding light arising from different sources, each with different intensities and existing on different timescales. This light is composed of the following components; the amplified spontaneous emission (ASE) pedestal, pre-pulses and uncompensated dispersion, the details of which will now be covered.

The ASE pedestal [82] arises during the amplification stages of the pulse generation process. Once population inversion is achieved in the gain medium, but before the arrival of the main pulse that triggers stimulated emission, a fraction of the excited electrons spontaneously decay to their ground state, emitting photons in a process known as spontaneous emission. These photons are emitted in random directions, with a percentage propagating along the central axis of the amplifiers. This leads to their amplification and thus the formation of the ASE pedestal. This feature can exist on

the timescale of tens of nanoseconds before the arrival of the main pulse, corresponding to the timescale on which the gating mechanisms in the amplification stages operate, hence the ASE pedestal cannot be completely removed. As spontaneous emission is an intrinsic process in any population-inverted medium, it cannot be eliminated entirely.

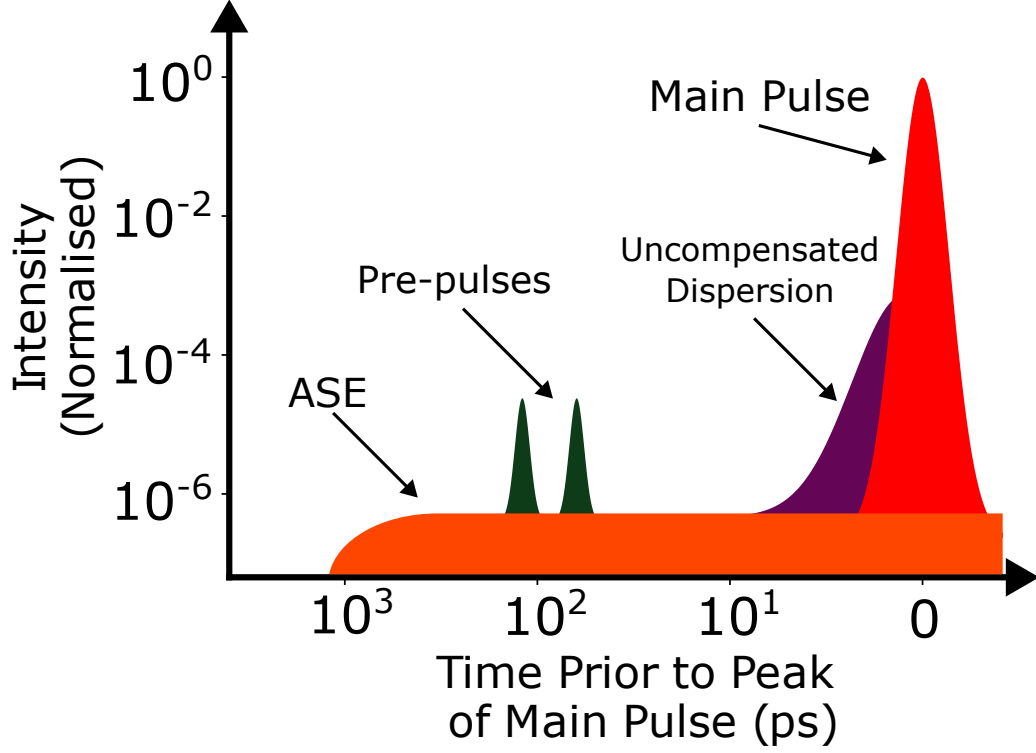


Figure 3.2: Illustration showing a simplified version of the temporal intensity profile of a typical pulse from a high-power laser system.

The pre-pulse component of the preceding light is a result of internal reflections within the laser-chain optics. As the stretched pulse passes through these optics, ‘post-pulses’ are created. The interaction between these pulses and the main pulse leads to spectral modulation of the main pulse. Compression of this spectrally modulated pulse in the final stage of the CPA process leads to the generation of pre-pulses. Such pulses typically exist on the timescale of picoseconds prior to the main pulse [83]. Pre-pulses can be suppressed through the use of wedge-optics. These are optics for which the two surfaces of the component are slightly angled rather than parallel. Internal reflections from these optics are directed away from the main pulse axis, such that they cannot interact with the main pulse, hence eliminating the pre-pulses [84]. Finally, uncompensated dispersion results from imperfections in pulse compression during the final stages

of the CPA process. This manifests on the rising edge of the main pulse, reducing the contrast on the picosecond timescale. Figure 3.2 shows an illustration of a typical temporal intensity profile of a high-intensity laser pulse.

Increasing the pulse contrast through modifications to the laser system itself, such as the use of wedge optics mentioned above, is not easily achievable. Instead, contrast enhancement techniques are often applied external to the pulse generation process but prior to the occurrence of the laser-plasma interaction being studied. The technique employed during the process of the experiments reported on in this thesis is the planar plasma mirror, which will now be discussed.

Plasma mirrors

A planar plasma mirror (PPM) is an optical component commonly used to increase the contrast of a pulse [85]. This typically consists of a planar piece of dielectric glass that is transparent to the laser wavelength. The operation of the PPM is as follows; it is positioned in the beam-line such that the lower intensity portion of the light preceding the pulse is below the ionisation threshold, allowing for it to be largely transmitted through the PPM. Complete transmission of this light does not occur as the reflectivity of the PPM to the laser wavelength is not 0%, although an anti-reflective (AR) coating is often applied to achieve reflectivities of $\sim 0.3\%$ [86]. After transmission of this lower intensity portion of the pulse, the rising edge of the main pulse begins to interact with the PM. As its intensity exceeds the ionisation threshold, a plasma is formed on its surface. This generated plasma remains overdense (section 2.4.3) for a period great enough that the remainder of the pulse, consisting largely of the main pulse, is reflected. However, the reflectivity of the overdense plasma is not 100%, as a portion of the pulse energy is absorbed into the plasma via the mechanisms discussed in section 2.5 of the previous chapter. This results in the energy present in the main pulse being reduced by between 15 and 30%, and thus a reduction in peak pulse intensity by the same value. In this way, the PM has enhanced the pulse contrast by up to two orders of magnitude [87], while the intensity of the main pulse has been largely maintained.

To ensure that the spatial qualities of the pulse are maintained after its reflection from the overdense plasma, the PM's surface must be sufficiently flat, often termed optical-quality. This means that deviations no larger than $\lambda_L/10$ are present across

its surface. This ensures that the overdense plasma that reflects the main pulse is of sufficient flatness that the wavefront or spatial properties of the main pulse are not altered during reflection. Figure 3.3(a) shows an illustration of the operating principle of a plasma mirror, with figure (b) showing a picture of a PPM taken from the experiment reported on in chapter 7 of this thesis.

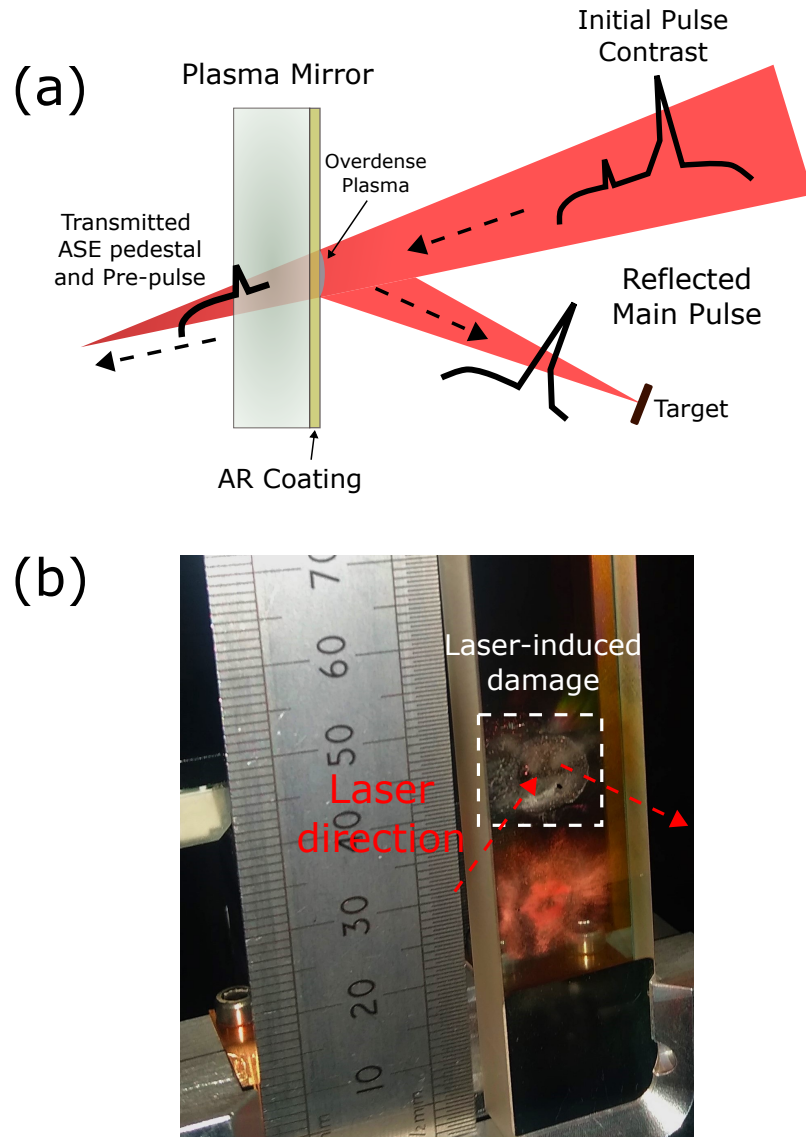


Figure 3.3: (a) Illustration showing the operating principle of a plasma mirror used to enhance the temporal intensity contrast of a high-power laser pulse. (b) An image of a planar plasma mirror that was utilised during an experiment performed using the Vulcan laser system, reported on in chapter 7.

3.3 High power laser facilities

Having outlined the general method for generating and amplifying pulses to relativistic intensities through the processes of CPA, this section provides an overview of the two high-power laser systems that were used to obtain the experimental data reported on in chapters 6 and 7 of this thesis, namely Vulcan and Gemini.

3.3.1 Gemini laser system

The results reported on in chapter 6 of this thesis are from an experimental campaign carried out using the Gemini laser system [88][89]. This facility consists of two sections, the first, ‘Astra’, is responsible for the generation and initial amplification of the pulse and the second, ‘Gemini’, is responsible for the final stage of amplification. Astra is a 10 TW system possessing a Ti:Sapphire gain medium, generating pulses with a central wavelength of 800 nm and a bandwidth of 40 nm. The Astra system generates a seed pulse that is then stretched from its initial duration of ~ 10 fs to 1.06 ns. The pulse then undergoes three stages of amplification, increasing its energy to 1.5 J. The diameter of the beam, and therefore amplifier optics, are increased over the successive amplification stages such that the pulse intensity does not exceed the damage threshold of the amplifier optics. The pulse is then directed to Gemini in which the final stage of pulse amplification takes place. The Gemini system is configured for dual beam capability, designed to allow for more exotic experimental setups such as beam-beam interactions [90]. This is enabled by splitting the pulse from Astra into two pulses with equal energies, both of which are then amplified. This final amplification stage increases the energy of both pulses to 15-20 J and the compression stage of the CPA process reduces its duration down to 40 fs. The Gemini system is capable of generating maximum energy pulses once every 20 seconds. With the use of an F/2 off-axis parabola, commonly employed during experiments conducted at this facility, intensities in the range of 1×10^{21} W/cm² are achievable. A typical pulse generated from this system will have a temporal intensity contrast on the order of 10^{10} at nanoseconds and 10^8 at picoseconds before the main pulse. The Gemini facility possesses a retractable dual plasma mirror system, the use of which can lead to pulse contrast enhancements of up to 10^4 [91]. The alignment of the pulse generated from Gemini is maintained through the use of a deformable mirror known as an adaptive optic (AO). This allows for the correction

of wavefront aberrations and pointing instability created through thermal effects that arise from continuous use of the system.

3.3.2 Vulcan laser system

The Vulcan laser system, also located at the CLF, is the second to be discussed in this thesis. This laser is currently shut down for several years to enable a substantial upgrade via the Vulcan 20-20 project [19] but at the time of this thesis work it supported two target-areas, target-area petawatt (TAP) and target-area West (TAW) [92], with the latter utilised for the experiment reported on in chapter 7 of this thesis. TAW operates with a total of eight beam-lines, consisting of two short-pulse (1-10 ps) and six long-pulse (0.8-5 ns). The two short-pulse beam-lines, referred to as beam 7 and beam 8, were used during the experiment reported on in this thesis and so only those will be discussed in this section.

Pulse generation within the Vulcan system begins with a seed pulse being produced from an oscillator containing a neodymium doped glass gain medium. This seed pulse has a central wavelength of 1053 nm and a bandwidth of 15 nm. Initially containing an energy of ~ 5 nJ, this pulse is then stretched to a duration of several nanoseconds before undergoing amplification. Similar to the Gemini system, the beam diameter is increased with each stage of amplification to prevent the damage threshold of the optics being exceeded. This necessitates a corresponding increase in gain medium diameter, beginning at 9 mm then ending at 150 mm in diameter. After undergoing six stages of amplification, the pulse is split into two pulses of equal energy to form beam 7 and beam 8. A final two stages of amplification increase the energy in both pulses to 100 J. Beam 8 can undergo one additional stage of amplification to deliver a pulse energy of 500 J but this capability was not utilised during the experiment. After the amplification is complete, the pulses are then compressed. Due to the relatively narrow bandwidth, compared to systems like Gemini, the shortest pulse duration achievable by Vulcan is approximately 1 ps. This system is capable of generating maximum energy pulses once every hour. During the experimental campaign reported on in this thesis the pulses delivered into TAW were brought to focus using an F/3 off-axis parabola, reaching peak intensities of $\sim 5 \times 10^{19}$ W/cm². Similar to the Gemini system, an AO is used to allow for the correction of wavefront aberrations.

3.4 Experimental diagnostics

This section will cover the diagnostic techniques used to characterise a laser-plasma interaction for the purpose of investigating the underlying physics. While there exists a large array of LPI specific diagnostics, it is not always practical or necessary to make use of many of these during a given experiment. Rather, it is the aspects of the interaction that are being studied that dictates the choice of diagnostic that are employed. For the results presented in this thesis, those utilised were optical diagnostics measuring the transmitted and generated light at the rear of targets and the technique of Stokes polarimetry, used to measure the polarisation state of light present at the rear of targets.

3.4.1 Transmitted and generated light

A key diagnostic used in the experiments reported on in chapters 6 and 7 of this thesis is that of measuring the quantity and spatial intensity distribution of the light at the rear of a target, arising from both the transmission of the main pulse and the generation of light in the plasma. This can provide important insights for key aspects of the interaction, such as the onset of transparency in expanding targets and the conversion of laser pulse energy into frequency-doubled light from aperture targets. This light is diagnosed by positioning a scatter screen to the rear of the target. The distance between this screen and the target is set such that the light can expand sufficiently so that its intensity is below the damage threshold of the screen. However, it cannot be placed such that the light over-fills the screen, as this would result in the measurement being an underestimate of the total signal present at the target rear. The specific distances from the target to the scatter screen are detailed in chapters 6 and 7, where the implementation of this diagnostic is discussed for each experiments. A scatter screen, typically composed of PTFE, acts as a Lambertian surface, diffusely scattering light that is incident upon it. Consequently, light that is scattered from such a surface has an equal intensity, independent of viewing angle.

The imaging of the light collected by the scatter screen is achieved through the use of CCD cameras equipped with telephoto lenses, which can be filtered for various frequencies when required. A schematic of this setup in which the ω_L and $2\omega_L$ frequencies are filtered for is shown in figure 3.4.

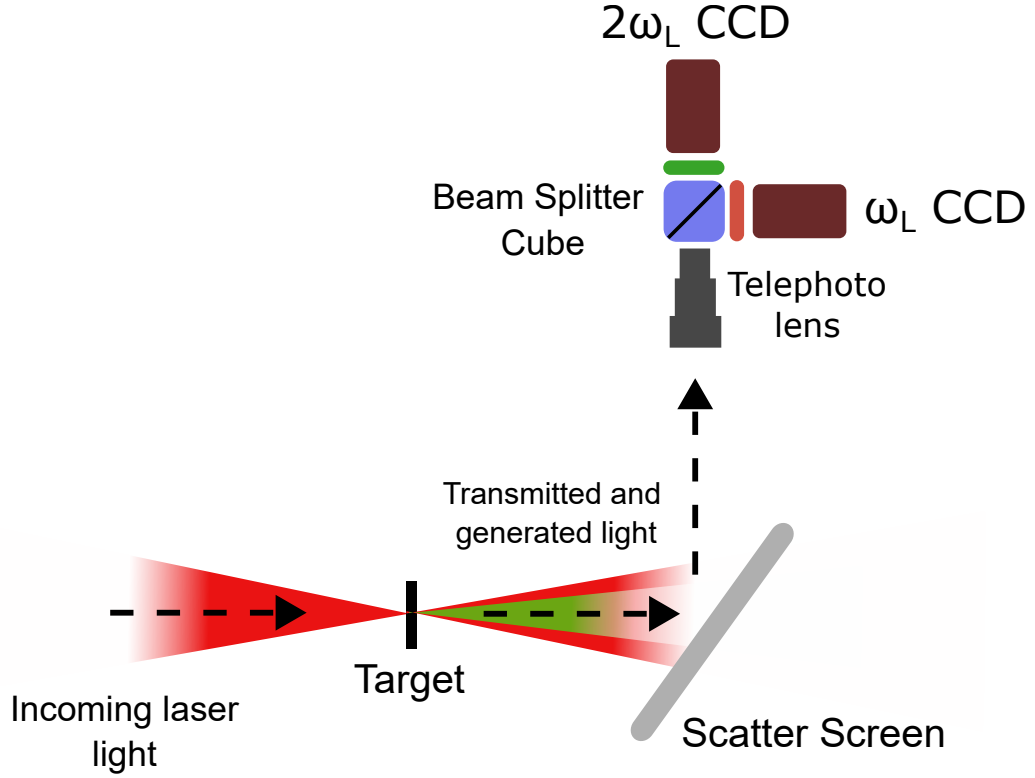


Figure 3.4: Illustration of a diagnostic setup utilised to measure the light transmitted through and generated in a target plasma. Light incident on the scatter screen is imaged by CCD cameras equipped with telephoto lenses, filtered for the fundamental and doubled frequencies.

An important consideration with respect to the CCD cameras used in this diagnostic is their bit depth, which refers to the number of discrete intensity levels that each pixel can represent. The CCD cameras typically used in this form of diagnostic have a bit depth of 16, allowing each pixel to represent up to 65,536 intensity levels. This figure is related to the cameras dynamic range, defined as the ratio of the maximum signal a pixel can detect before saturating to the minimum detectable signal. Despite a 16-bit CCD camera possessing a dynamic range far greater than a typical consumer camera, it is still unable to capture the full range of signals present in the light collected on the scatter screen in this diagnostic. As the intensity of this light typically spans many orders of magnitude, direct imaging without any attenuation would very likely saturate the camera, resulting in the loss of information from the brightest regions. To address this, filtering is applied so that the intensity of light from typical shot produces a signal that lies near the center of the dynamic range of the camera. The level of filtering can initially be determined through calibration shots, typically performed without a target.

Such shots often use pulse energies below the maximum that is available, chosen to approximate the intensity of light expected during full-power shots for which a target is present. As the experiment enters the data collection stage, initial target shots may reveal that the filtering requires further adjustment such that the desired signal level is achieved. Further details of the calibration shots taken during the experiments are covered in their respective chapters. The form of filtering employed to prevent saturation of the CCD cameras is termed a neutral density (ND) filter. The intensity of any light that propagates through an ND filter is attenuated, with the degree of attenuation indicated by its ‘ND number’. A higher value indicates a greater level of attenuation.

Calibrations of the CCD cameras are also performed to characterise their response. This involves exposing a camera to light of increasing intensity and recording the maximum signal the camera detects at each level. The light intensity is increased until saturation of the camera occurs. The resulting relationship between the measured signal and the incident light intensity is then plotted. This can reveal any non-linearities in the camera’s response that can subsequently be corrected during the analysis of measurements made with the camera.

3.4.2 Spectral measurements

The spectral intensity of light was measured during the two experiments using optical spectrometers, capable of detecting wavelengths between 400 and 1100 nm. These were employed both at the rear of a target and as part of a Stokes polarimeter diagnostic. Details of the latter are discussed in the next section of this chapter and in the ‘diagnostics’ section of chapter 7. The principle of operation of the spectrometers used is as follows. The light being characterised enters the spectrometer and is incident on a diffraction grating, spatially separating the light by wavelength. This dispersed light is then directed onto a CCD detector that measures the intensity of the various wavelengths of light. As with the CCD cameras discussed in the previous section, ND filters are employed with the spectrometers to prevent saturation of the detector. The appropriate level of ND filter is determined using the same procedure as for the CCD cameras.

Spectrometer calibration is performed to ensure the correct response of the CCD detector to a range of wavelengths of light. This is achieved using a light source with

a well-characterised spectral intensity profile, against which the spectrometer output could be compared. Any deviations from this known spectral profile can then be accounted for in subsequent analysis of the measurements made by the spectrometer.

3.4.3 Stokes polarimetry

The results presented in chapter 7 of this thesis involve the measurement of the polarisation state of light present at the rear of ultra-thin targets using the technique of polarimetry. The specific implementation of polarimetry that was used was that of Stokes polarimetry [93]. This technique characterises the polarisation state of light through a Stokes vector, each component of which is referred to as a Stokes parameter. These parameters can be measured experimentally and are related to the total intensity of the light passing into the polarimeter diagnostic, as well as the intensity of the six so-called ‘degenerate states’ commonly employed in polarimetry. These parameters are defined by the following equations.

$$S_0 = I_0 + I_{90} = I = 1 \quad (3.4)$$

$$S_1 = I_0 - I_{90} = Ip \cos 2\psi \cos 2\chi \quad (3.5)$$

$$S_2 = I_{45} - I_{135} = Ip \sin 2\psi \cos 2\chi \quad (3.6)$$

$$S_3 = I_{RCP} - I_{LCP} = Ip \sin 2\chi \quad (3.7)$$

where I_n is the intensity of light of degenerate state n , I is the total intensity of the light passing into the polarimeter diagnostic and p is the degree of polarisation, taking values of $0 \leq p \leq 1$. The variables 2ψ and 2χ are so-called characteristic angles, discussed further below. The degenerate polarisation states refer to linear 0° , 90° , 45° , and 135° , and right and left handed circular polarisation. The linear angle of polarisation in this case is defined relative to the plane orthogonal to the light propagation direction. The assumption is often made in the context of plasma polarimetry, that the degree of polarisation of the light is 1, hence stating that the S_0 parameter is equal to 1 in the equations stated above.

While the specific way in which this diagnostic was implemented during an experiment is covered in chapter 7, a typical Stokes polarimetry diagnostic involves splitting the light under investigation into three separate paths. Each path contains an optical component, such as a polariser, that allows for the orthogonal polarisation states to be separated. The intensity of the light of each of these channels is then recorded and used to calculate the individual Stokes parameters.

Stokes vectors can be visualised using the so-called ‘Poincaré sphere’ in which a polarisation state represented by a given Stokes vector is indicated by the direction and magnitude of a vector in 3D space. Figure 3.5 shows an example. Each axis of the sphere represents one of the Stokes parameters. For example, the Stokes vector representing light that is purely linearly polarised along the horizontal axis,

$$\begin{bmatrix} 1 \\ 0 \\ 1 \\ 0 \end{bmatrix}$$

would be represented by a vector pointing along the S_2 axis, ending at the coordinates (1,0,0). Any change in the linear polarisation state is then represented by a movement of the vector along the horizontal plane of the sphere, toward the S_1 axis. Any change in the polarisation state from purely linear to elliptically polarised is represented by a movement of the vector away from the horizontal axis of the sphere, along its latitude, indicating an increase in the magnitude of the S_3 component. The magnitude of a vector is determined by the degree of polarisation. For fully polarised light, the tip of the vector will end on the surface of the sphere, while for partially polarised light, the tip will lying inside it.

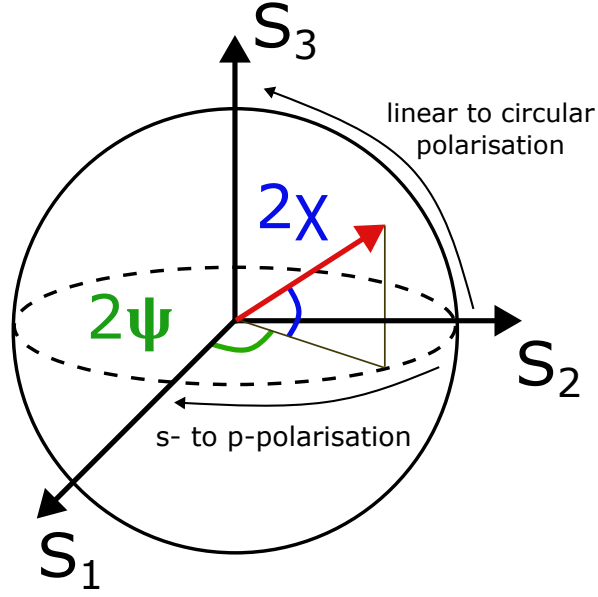


Figure 3.5: Illustration of the Poincaré sphere, used to visualise the Stokes vector representation of polarised light.

The characteristic angles used to define the Stokes parameters can be seen to correspond to the angles the polarisation vector makes relative to the horizontal and vertical axes of the Poincaré sphere.

3.5 Numerical simulations

With the advent of high-performance computing (HPC), it has become standard practice to complement experimental measurements with numerical simulations. These offer insights into the dynamics of interactions that cannot be captured through experimental diagnostics alone, making them indispensable for understanding the underlying physical processes. For example, the duration of certain interactions, typically on the order of femtoseconds to picoseconds, falls several orders of magnitude shorter than the response times (nanoseconds) of the components of diagnostics that make measurements, such as the CCD cameras. Consequently, the majority of experimental measurements are time-integrated, which limits the ability to resolve intra-pulse dynamics through direct observation. In addition, despite the extensive array of LPI diagnostics available, the limited number of variables that can be measured experimentally constrains the depth of understanding achievable regarding the physics under investigation. In contrast, numerical simulations can model interactions on sub-femtosecond timescales

and provide data on a significantly larger set of variables. This capability allows for a more comprehensive understanding of the complex physical phenomena being studied. It is not practically feasible to fully simulate an entire interaction, so different forms of numerical simulations are employed based on the specific physical processes most relevant to the aspect of the interaction under study. The various forms of numerical simulation available provide a simplified representation of a particular portion of the interaction.

In this thesis, two simulation types were used: particle-in-cell (PIC) and hydrodynamic. As discussed in the following sections, these methods are tailored to address different aspects of the interaction, each requiring different algorithms, numerical schemes, and computational resources to achieve accurate results.

3.5.1 Particle-in-cell simulations

Particle-in-cell (PIC) codes are used to model interactions occurring on femtosecond to picosecond timescales under high- to ultra-high-intensity laser pulses, and are typically the most computationally demanding form of simulation for laser-plasma interactions. The simulation of such interactions often requires the computational load to be distributed across multiple CPUs in a HPC cluster, with parallel processing employed to manage the large number of calculations. The memory requirements for such simulations can reach into the hundreds of gigabytes of RAM. Furthermore, the volume of data generated by PIC simulations often scales to several gigabytes, far exceeding the storage capacity of a typical desktop computer.

The principle behind the use of PIC codes involves describing the plasma using a so-called kinetic model. A key assumption of this approach to modelling a laser-plasma interaction is the fact that the system being modelled is not in local thermal equilibrium, an assumption that distinguishes this approach from hydrodynamic simulations. A 6-dimensional distribution function is employed that represents both the momentum and position of a particle in three dimensions. The fully kinetic evolution of the distribution function includes the effects of both electromagnetic and collisional physics on the particle, but for the purposes of modelling the interaction between the highest intensity portion of the laser-pulse and the plasma, collisional physics does not have a meaningful impact. As such, only the electromagnetic forces are considered, and the evolution of the plasma can be found by solving the Vlasov equation:

$$\frac{\partial f}{\partial t} + \mathbf{v} \cdot \frac{\partial f}{\partial \mathbf{x}} + q(\mathbf{E} + \mathbf{v} \times \mathbf{B}) \cdot \frac{\partial f}{\partial \mathbf{p}} = 0 \quad (3.8)$$

This is a form of the Fokker-Planck equation in which collisional effects, which could be represented by a term on the right-hand-side of the equation, are ignored. While the removal of the collisional component of the interaction aids in reducing the complexity of equation 3.8, given that the electron number density of a typical material used in laser-plasma interactions is on the order of 10^{27} m^{-3} , directly solving the Vlasov equation with this number of particles is not practically possible. To reduce the computational resources required, the PIC method uses a much lower number of what are termed ‘macroparticles’. A single macroparticle can represent a large number of real particles, leading to a reduction in the computational resources required to solve the Vlasov equation, putting it in line with the HPC systems. This is a valid simplification as the macroparticles respond to electric and magnetic fields in the same way as the particles they represent, thus the behaviour of the plasma throughout the interaction is accurate.

An iterative process is employed to calculate the evolution of a plasma throughout the interaction. This begins with the macroparticles being mapped onto a fixed grid, with this mapping initially being specified by the user. The density and current due to this mapping are then calculated. These quantities are then used to calculate the resultant $\mathbf{E}$ and $\mathbf{B}$ fields using Maxwell’s equations. These fields, in combination with Lorentz force equation, are then used to calculate the motion of the macroparticles, leading to updated density and current values, and so the next iteration of the loop begins. This process is illustrated in figure 3.6.

As invaluable as PIC simulations have become as an aid to understanding laser-plasma interactions, they are not without their limitations and as such care must be taken to avoid commonly encountered issues. The first of these potential issues is that there must be a sufficient number of macroparticles per grid cell, as this value determines the ‘smoothness’ of the distribution function and the level of statistical noise in the simulation outputs. Signatures of a given physical process in the data could be masked or drowned out by a large amount of background noise, leading to insights being missed.

Spatially, the grid resolution must be larger than the Debye length. If this condition is not met the plasma temperature will increase until it can be resolved. This process, known as numerical heating, can lead to non-physical results and as such must be

avoided. To ensure that the issues detailed above are not encountered, an iterative process known as convergence testing can be performed in which a simulation parameter, macroparticles-per-cell for example, is increased and simulation outputs are monitored. Once a change in the simulation input has no effect on the output, the ‘correct’ PPC number has been converged upon.

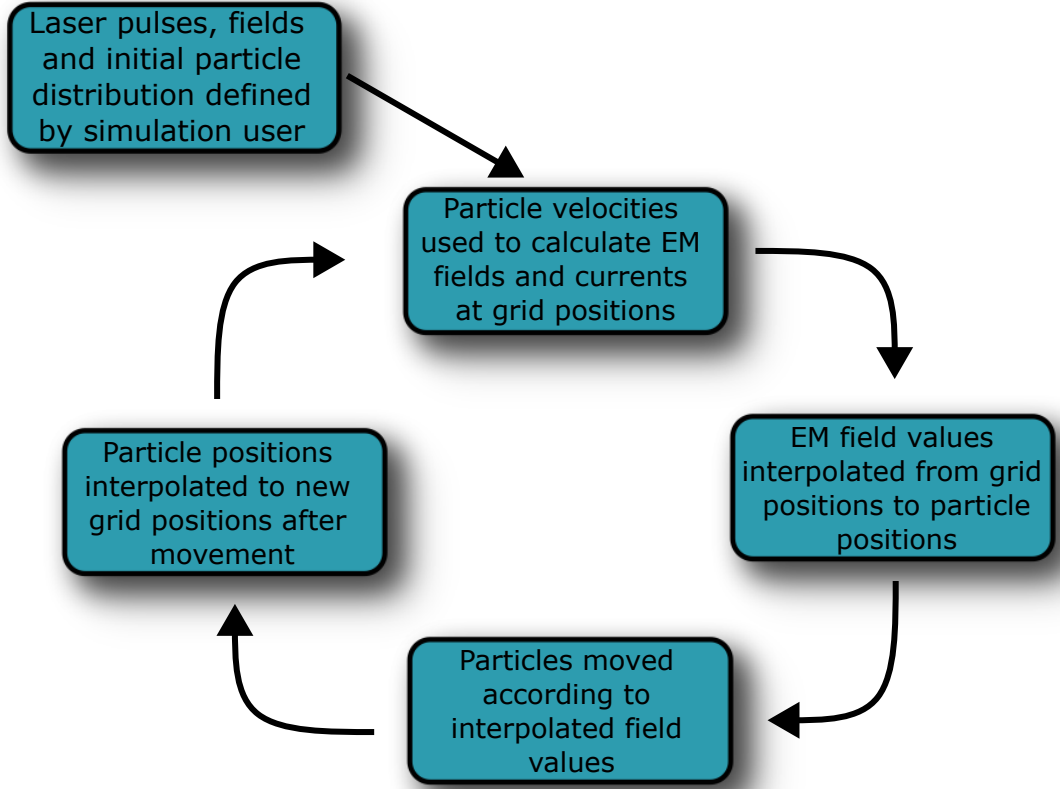


Figure 3.6: Flow diagram displaying the PIC algorithm.

The specific PIC code used for the work presented in this thesis was EPOCH [94]. This code is capable of running simulations in 1- 2- or 3-dimensions and is distinguished from many other available PIC codes by being open source. This enables users to edit the simulation code to suit their needs. This code employs a message passing interface (MPI) protocol, which distributes the necessary calculations between a large number of CPU’s, reducing the time required to perform a simulation. In addition, this code is designed to be user-friendly, with simulations parameters such as simulation domain, laser intensity and initial target particle distribution, easily modified via an input deck file. The specific input parameters used for the simulations reported on in this thesis are provided in the relevant chapter sections.

Other widely used open-source PIC codes include Smilei and WarpX, which offer complementary capabilities. Smilei is capable of modelling relativistic laser-plasma simulations with a strong emphasis on scalability and modular physics packages, using high-performance computing [95], while WarpX targets extreme-scale computing architectures and advanced electromagnetic solvers for plasma-based accelerator modelling [96].

3.5.2 Hydrodynamic simulations

Hydrodynamic codes take a fundamentally different approach to modelling laser-plasma interactions. With the assumption that the system is in local thermal equilibrium, individual particle distributions are neglected and the plasma is treated as a fluid. Bulk properties of the plasma are simulated, such as pressure and fluid velocity, making them well-suited for modelling interactions on timescales of hundreds of picoseconds to nanoseconds. Not suitable for simulating high-intensity pulses, this form of simulation is often used to gain an understanding of the scale-length of the pre-plasma that is generated by a given laser system's pre-pulse and ASE pedestal, as these features often exist on the pico- and nanosecond time scale and at lower intensities. In these simulations, the user specifies the initial spatial distribution and values of plasma properties, such as pressure, density, and velocity, which are mapped onto a spatial grid, along with the characteristics of the laser used in the interaction. Unlike in particle-in-cell (PIC) simulations, where the laser pulse is modelled through electric and magnetic fields governed by Maxwell's equations, hydrodynamic simulations generally treat the laser as a simple source of energy input, with the way in which the energy is absorbed by the plasma varying between the specific modelling code in use. Once the initial conditions are set, the hydrodynamic equations are solved, updating the plasma's properties across the spatial grid. A unique aspect of hydrodynamic simulations is the handling of the spatial grid itself. Several coordinate systems are used, such as the Eulerian system, where the grid is fixed and the fluid moves relative to it, and the Lagrangian system, where the grid moves with the fluid. Additionally, Arbitrary Lagrangian-Eulerian (ALE) approaches exist, combining both systems and switching between them as needed.

The hydrodynamic simulations presented in this thesis were conducted using the FLASH code, developed initially by the University of Chicago and later by the University of Rochester. FLASH employs the Eulerian approach and features adaptive

mesh refinement (AMR), which adjusts the grid resolution based on computational requirements, optimising resource use and reducing runtime. Laser energy deposition in FLASH is modelled using inverse Bremsstrahlung absorption. The specific input parameters and convergence testing of the FLASH code is covered in detail in chapter 5 of this thesis.

3.6 Summary

In summary, this chapter provides an outline of the experimental and simulation-based methodologies used to gain an understanding of novel aspects of laser-plasma interaction physics presented in this thesis. Beginning with the outline of a typical approach to conducting experimental campaigns in this field, the key technologies involved in generating high-intensity laser pulses are explored. The challenges associated with their use and the mitigation strategies employed, namely plasma mirror technology, are then discussed. Subsequently, the chapter provides an overview of the two high-power laser facilities utilised in the experimental campaigns discussed in this thesis, Gemini and Vulcan. Following this, the diagnostic tools used during these campaigns are described. Finally, the chapter discusses the numerical simulation methods, those being particle-in-cell (PIC) and hydrodynamic simulations, that were employed to aid in the understanding of the experimental measurements.

Chapter 4

Numerical investigation of higher order modes of intense second harmonic light generated from a plasma aperture

4.1 Introduction

This chapter reports on the investigation of the interaction between relativistically intense laser pulses and aperture targets that can generate frequency-doubled light with higher-order spatial mode structures. Through a combination of 2D and 3D EPOCH simulations, the study identifies the interaction parameters that maximise the conversion efficiency of energy from the laser pulse into $2\omega_L$ light. Key factors were found to include the fraction of light transmitted through the aperture, the strength of the longitudinal electric field at the aperture edge, and the absorption of laser pulse energy as it propagates through the target. The investigation extends to 3D simulations, allowing for the investigation of the effects of varying the drive pulse polarisation from linear to circular. This change is found to result in an increase in the efficiency of $2\omega_L$ generation and a change in the generated mode structure from a TEM_{01} mode to an LG_{01} , prompted by the polarisation vector being perpendicular to the aperture edge throughout the interaction in the case of circular polarisation. Additionally, the impact of pulse temporal intensity contrast and alternative aperture geometries on conversion

efficiency is investigated. This resulted in a decrease in efficiency of $2\omega_L$ generation with lower contrast but the ability to increase this efficiency by tailoring the target geometry to the focusing characteristics of the pulse. These findings demonstrate the potential of preformed aperture targets as plasma optical devices for mode conversion of intense light, providing a foundation for subsequent experimental investigations.

4.1.1 Light generation from a plasma aperture

The investigation reported in this chapter builds upon prior work that demonstrated the formation of a self-generated relativistic plasma aperture during the interaction of an intense laser pulse with an ultra-thin (tens of nm) target [56]. In these interactions, an initially opaque region of the target becomes relativistically transparent to the laser light, allowing for transmission and diffraction of the drive pulse and manipulation of the particle dynamics in the corresponding field structures [61][57].

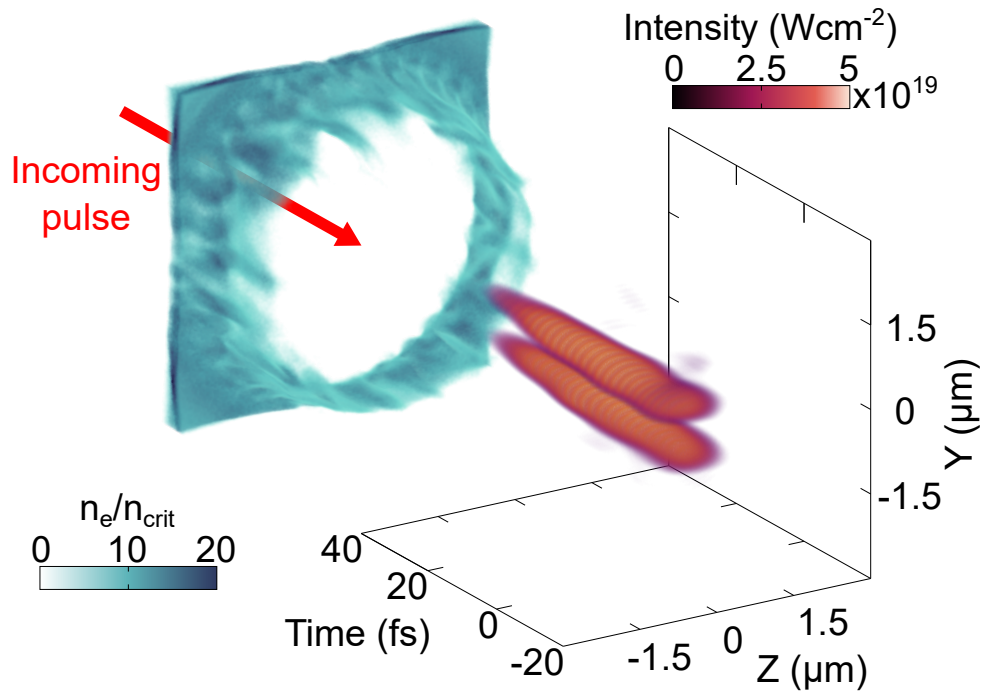


Figure 4.1: Example 3D simulation results illustrating the temporal behaviour of relativistically intense frequency-doubled light in a higher-order spatial mode (TEM_{01}) at a distance of $5 \mu\text{m}$ from the rear of the target, driven by intense laser light in a fundamental TEM_{00} mode. The results depict a self-generated relativistic plasma aperture in an initially solid-density, 10 nm-thick aluminium target. Taken directly from [97].

The spatial profile of the incident laser pulse determines the shape of this transparent region, which, for a Gaussian focal spot, forms an aperture with an approximately

circular profile. As the laser pulse propagates through this relativistically under-dense region, it interacts with the edges of the plasma aperture, generating intense light at both the laser frequency, ω_L , and higher harmonics. These harmonics exhibit higher-order transverse spatial modes and polarisation states distinct from the drive laser pulse. Figure 4.1 presents an example result from a 3D EPOCH simulation of a relativistically intense laser pulse interacting with an initially solid-density, 10 nm-thick aluminium target [56]. The simulation reveals the formation of a self-generated relativistic plasma aperture and the emission of $2\omega_L$ light in the TEM₀₁ mode.

For this and all subsequent figures presenting simulation-derived intensity values, it is the instantaneous rather than the cycle-averaged intensity that is shown, calculated using the following expression

$$I = \frac{\epsilon_0 c E^2}{2} \quad (4.1)$$

Since this calculation is used consistently throughout the thesis, it provides a representative measure of the generated light intensity across all results.

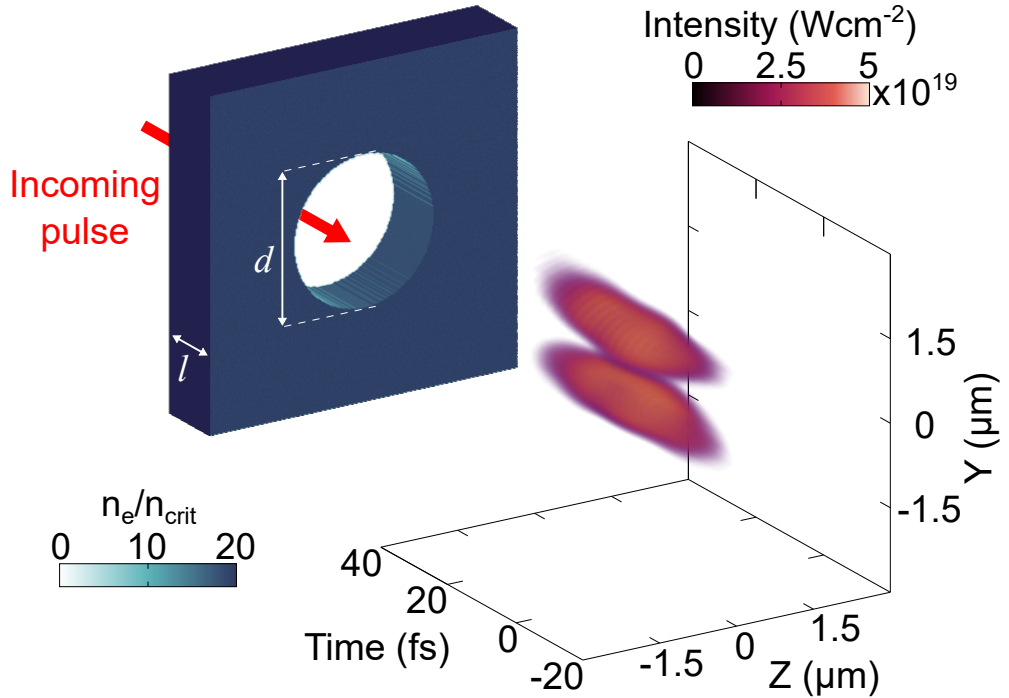


Figure 4.2: Example 3D simulation results illustrating the temporal behaviour of relativistically intense frequency-doubled light in a higher-order spatial mode (TEM₀₁) at a distance of 5 μm from the rear of the target, driven by intense laser light in a fundamental TEM₀₀ mode. The results depict a preformed aperture target with a thickness of 2 μm and an aperture diameter of 3 μm. Taken directly from [97].

This laser-plasma aperture interaction effectively transfers energy from the incident laser pulse, which has a fundamental Gaussian spatial mode structure, to intense harmonic light with higher-order spatial modes. While ultra-thin targets pose challenges due to their fragility which can impose limitations on the achievable shot repetition rate, the concept could be simplified by employing micrometer-thick targets with preformed circular apertures to achieve similar effects.

Figure 4.2 illustrates a 3D simulation of a relativistically intense pulse interacting with a preformed aperture target with a thickness of $2\ \mu\text{m}$ and an aperture diameter of $3\ \mu\text{m}$. As can be seen, light of a similar intensity and transverse spatial structure to that generated in the case of an ultra-thin target is produced. This demonstrates the feasibility of generating relativistically intense $2\omega_L$ light in the TEM_{01} spatial mode from pre-formed aperture targets.

4.2 2D Simulations

The generation of frequency-doubled light from an aperture target described above was initially studied through the use of 2D EPOCH simulations. Although mode generation has previously been investigated using 3D simulations, it can still be effectively studied in only 2 dimensions. The interaction between a linearly polarised laser pulse and the aperture edges that is responsible for the generation of $2\omega_L$ light occurs primarily in the X - Y plane, where the mode is also observable, allowing for meaningful analysis to take place. For this reason, only linearly polarised light in the Y direction was considered, as circular polarisation would not interact with the aperture edge in its entirety. However, limitations inherent to 2D simulations, such as the confinement of electron propagation to two dimensions, may lead to overestimations in some trends being investigated, such as the quantity of generated $2\omega_L$ light with varying pulse energy [98]. Nevertheless, the lower computational demands of 2D simulations allowed for a wider parameter space to be explored, with trends identified in 2D serving as a foundation for subsequent 3D simulations, which can then increase confidence in these findings.

The simulation grid was defined as 1000×720 mesh cells, corresponding to $20\ \mu\text{m} \times 20\ \mu\text{m}$. The targets were comprised of a slab of Al^{13+} ions of thickness l with an aperture of diameter d . This was neutralised with an electron population with a peak density of $30n_{\text{crit}}$ and an initial temperature of 10 keV. To reduce the computational

cost, the simulations were performed with a reduced peak density, (in comparison to the solid density of $\sim 430n_{crit}$). Each species was initialized with an average of 80 particles per cell. All the boundaries of the simulation box were defined as free-space, and the ions were mobile in all cases. The pulse of laser light with wavelength $\lambda_l = 800$ nm entered from the left (X) boundary and was focused at the plane corresponding to the front surface of the target ($X = 0$). Both the laser pulse temporal and spatial profiles were defined in terms of intensity as Gaussian with pulse duration of $t_L = 40$ fs (FWHM) and focal spot diameter of $\phi_L = 2.1 \mu\text{m}$ (FWHM). Time $t = 0$ was defined as when the peak of the pulse reaches $X = 0$. The peak laser pulse intensity was varied from $I_L = 1 \times 10^{21}$ to 1×10^{23} W/cm² by increasing the energy in the pulse. Figure 4.3 shows an example of the electron density distribution at a simulation time of $t = 80$ fs taken from a simulation in which a pulse with $I_L = 1 \times 10^{21}$ W/cm² interacts with an aperture with $d = 3 \mu\text{m}$ and $l = 1.5 \mu\text{m}$. Clearly seen are the electron bunches that have been accelerated as a result of the interaction that contribute to the generation of the frequency doubled light, discussed in a later section of this chapter.

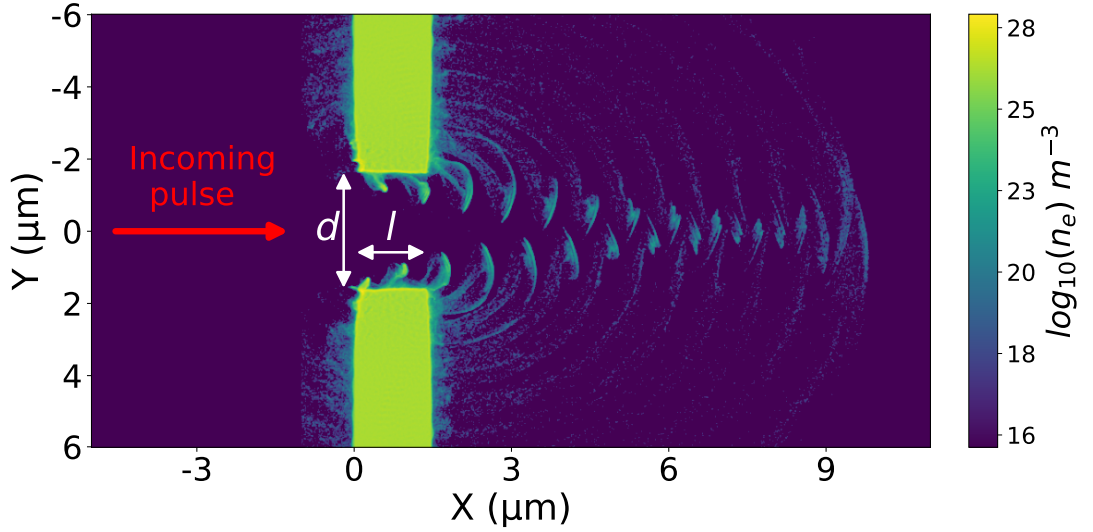


Figure 4.3: Electron number density taken from a 2D simulation in which a pulse with $I_L = 1 \times 10^{21}$ W/cm² interacts with a target of $d = 3 \mu\text{m}$ and $l = 1.5 \mu\text{m}$, at a simulation time of 80 fs. Labelled are the target thickness l and the aperture diameter d , along with the direction of propagation of the laser pulse.

4.2.1 Filtering simulation data for ω_L and $2\omega_L$ light

To investigate the interaction parameters influencing the generation of light from aperture targets, a method was developed to filter the simulation data for both the funda-

mental laser frequency and the doubled frequency. The process begins with applying a Fourier transform to the field data in the spatial domain, converting it into the frequency domain. A Gaussian bandpass filter is then applied to this transformed data. In this form of filtering, all signal outside the desired frequency range are removed, leaving only the desired frequency. After filtering, an inverse Fourier transform is performed, converting the E_Y field data back into the spatial domain. This frequency filtering technique was used to extract information about the generated $2\omega_L$ light, such as total energy and spatial distribution, from the simulation data.

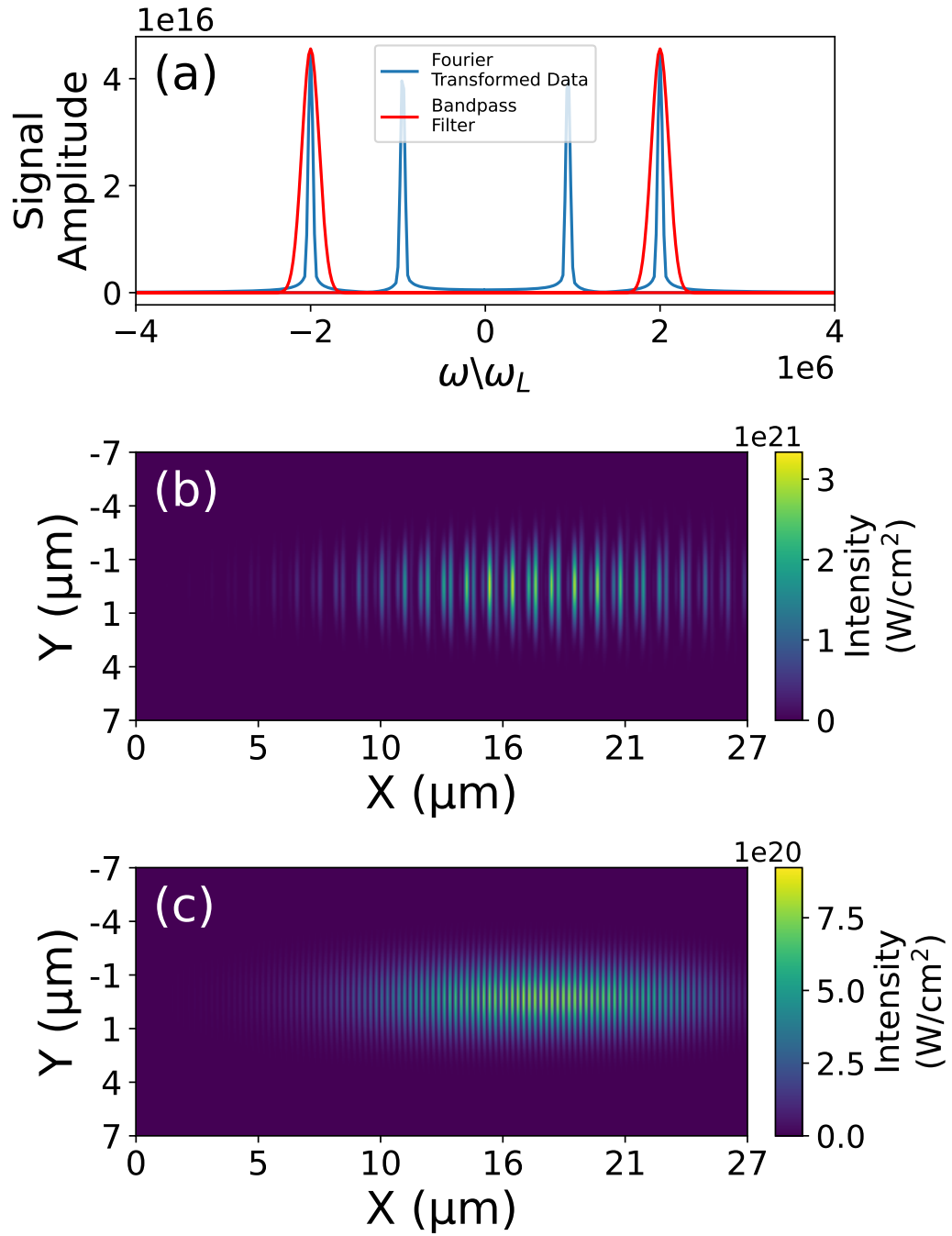


Figure 4.4: (a) The amplitudes of the signals representing the various frequencies of light generated from the aperture mechanism after the application of the frequency filtering method. Also included in this plot is the Gaussian bandpass filter applied to select for the frequency-doubled light. (b) and (c) show the spatial intensity distribution of the light present in the simulation data before and after the application of the frequency filtering method, respectively.

To validate the accuracy of this method, a 2D simulation was performed in which two overlapping pulses were propagated in a vacuum: one with the fundamental laser frequency and another with a doubled frequency. If the frequency filtering operates

correctly, it should allow for the selection of either frequency from this data. For verification, data from a single time step was chosen, and the method was applied to this snapshot. Figure 4.4(a) displays the field data transformed into the frequency domain, along with a Gaussian bandpass filter (red) centred to isolate the doubled frequency component. As seen, the two pulse frequencies are clearly distinguishable. Figures 4.4(b) and (c) present the original pulses and the filtered doubled-frequency pulse, respectively, confirming the filter’s effectiveness in isolating the desired frequency.

4.2.2 Investigation of the generation mechanism of frequency-doubled light

To begin the investigation of the use of pre-formed aperture targets for generating frequency-doubled light with higher-order spatial modes, interaction dynamics thought to contribute to $2\omega_L$ generation were analysed. These were the generation of electron bunches and their interaction with the longitudinal (E_X) electric field present at the rear of the target and within the aperture channel. Prior work presenting the generation of $2\omega_L$ light from a self-generated plasma aperture that this study builds upon [56] proposed that this light is produced from the constructive interference between the broadband emission from electron bunches being decelerated in the sheath field at the rear of an ultra-thin target. Given the similarities between pre-formed and self-formed plasma apertures, it is hypothesised that a similar mechanism is responsible for $2\omega_L$ generation in this context.

Displayed in figure 4.5(a) is the electron number density, normalised to the critical density, from a 2D simulation at time $t = -4.8$ fs in which a drive laser pulse with $I_L = 1 \times 10^{21}$ W/cm² interacts with an aperture target with $l = 0.5$ μm and $d = 3$ μm . Clearly seen is the emission of electron bunches from both edges of the aperture. As electron bunch acceleration is due to the oscillation of the lasers’ electric field vector, they are emitted from a single edge with a frequency of ω_L , π out of phase with those from the opposite edge, and with an overall frequency of $2\omega_L$ when considering both edges.

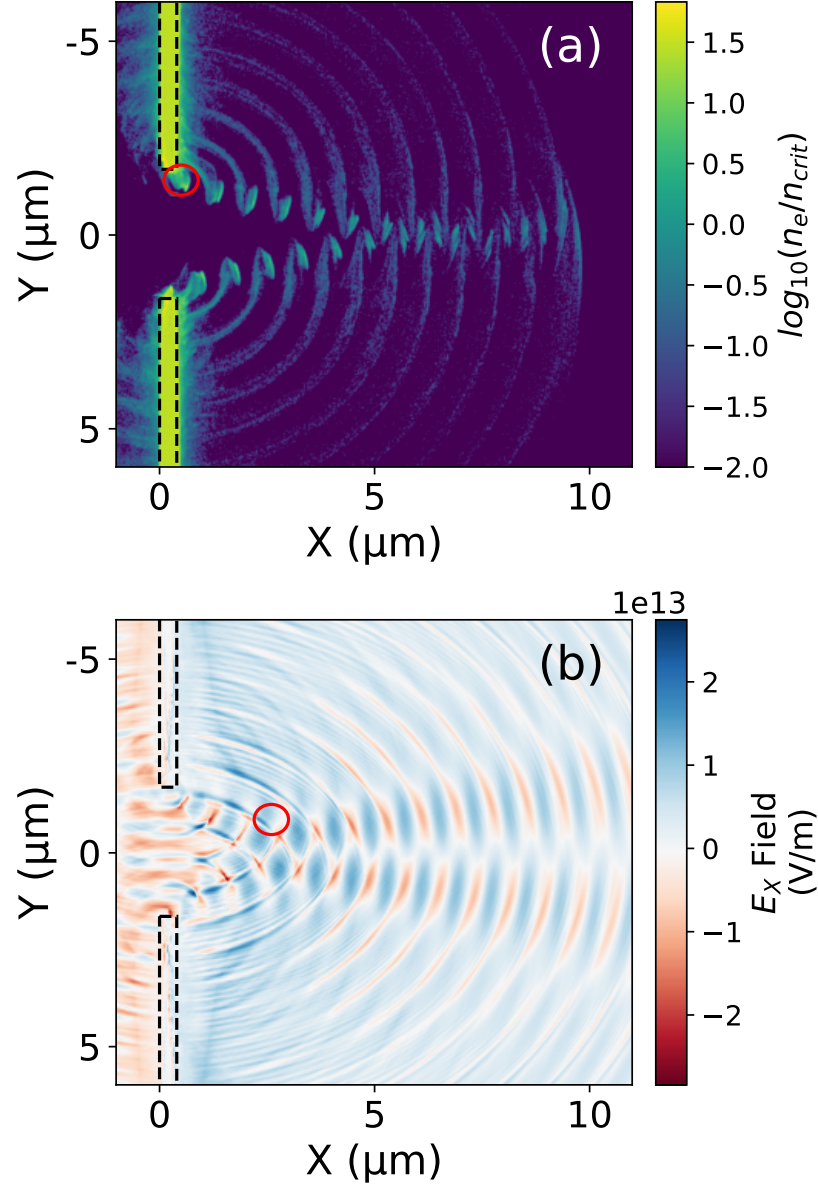


Figure 4.5: 2D simulation results for a target with $l = 0.5 \mu\text{m}$ and $d = 3 \mu\text{m}$: (a) electron density at $t = -4.8$ fs; (b) longitudinal electric field at $t = 2.8$ fs. The dashed lines indicate the initial positions of the target surfaces and aperture edges. An arbitrary electron bunch is circled in red for the purpose of tracking interaction metrics over time.

Figure 4.5(b) shows the E_X field from the same simulation at a later time of $t = 2.8$ fs. The interference of the fields emitted from the two opposite edges of the aperture, essentially point sources, creates the structure in the resulting fields seen in this 2D map. An arbitrary electron bunch, circled in red in figure 4.5(a) and again in figure 4.5(b) 6 fs later, is tracked throughout the simulation to allow for the investigation of the temporal variation in generated $2\omega_L$ light. Figure 4.6 shows the total kinetic energy K_e present in the bunch, plotted as a function of the axis of laser propagation,

X . In addition to this, the longitudinal electric field experienced by the bunch, E_X , and the intensity of the $2\omega_L$ light, $I_{2\omega_L}$ (both averaged over a length of 400 nm in the X -direction, centred on the electron bunch) are plotted. Note that arbitrary units are used to enable direct comparison of the field and intensity variation.

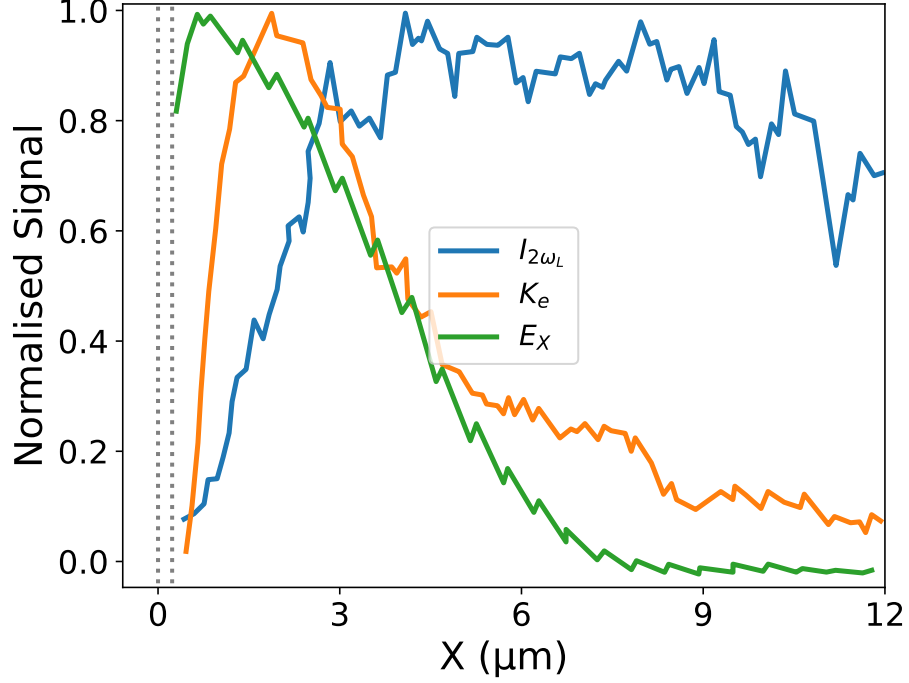


Figure 4.6: Total energy (K_e , orange) in the electron bunch marked by the red circle in figures 4.5(a) and (b) as it propagates in space, together with the intensity of the $2\omega_L$ light generated, averaged 400 nm around the bunch ($I_{2\omega_L}$, blue), and the longitudinal electric field experienced by the bunch (E_X , green) averaged over the same area. Note that arbitrary units are used for the quantities to allow for direct comparison.

Once emitted from the aperture edge, the bunch is rapidly accelerated to relativistic energies, after which the E_X field present at the rear of the target causes it to decelerate. This decelerating field exists due to the charge separation formed by fast electrons leaving the target prior to bunch generation and the rear surface plasma surrounding the aperture [99]. Evidence for this field can be observed in figure 4.5(b), in the positive electric field that is present for $X > 0$ (the rear of the target), along with the longitudinal field associated with the laser pulse. The generation of $2\omega_L$ light is then explainable by the interference between the broadband emissions from individual bunches from either aperture edge as they are accelerated and decelerated in the longitudinal electric fields. This leads to the intensity of the $2\omega_L$ light along the longitudinal axis increasing with distance from the aperture for $\sim 5 \mu\text{m}$, after which it saturates, then decreases as the

$2\omega_L$ light expands away from the sampled region in the vicinity of the bunch.

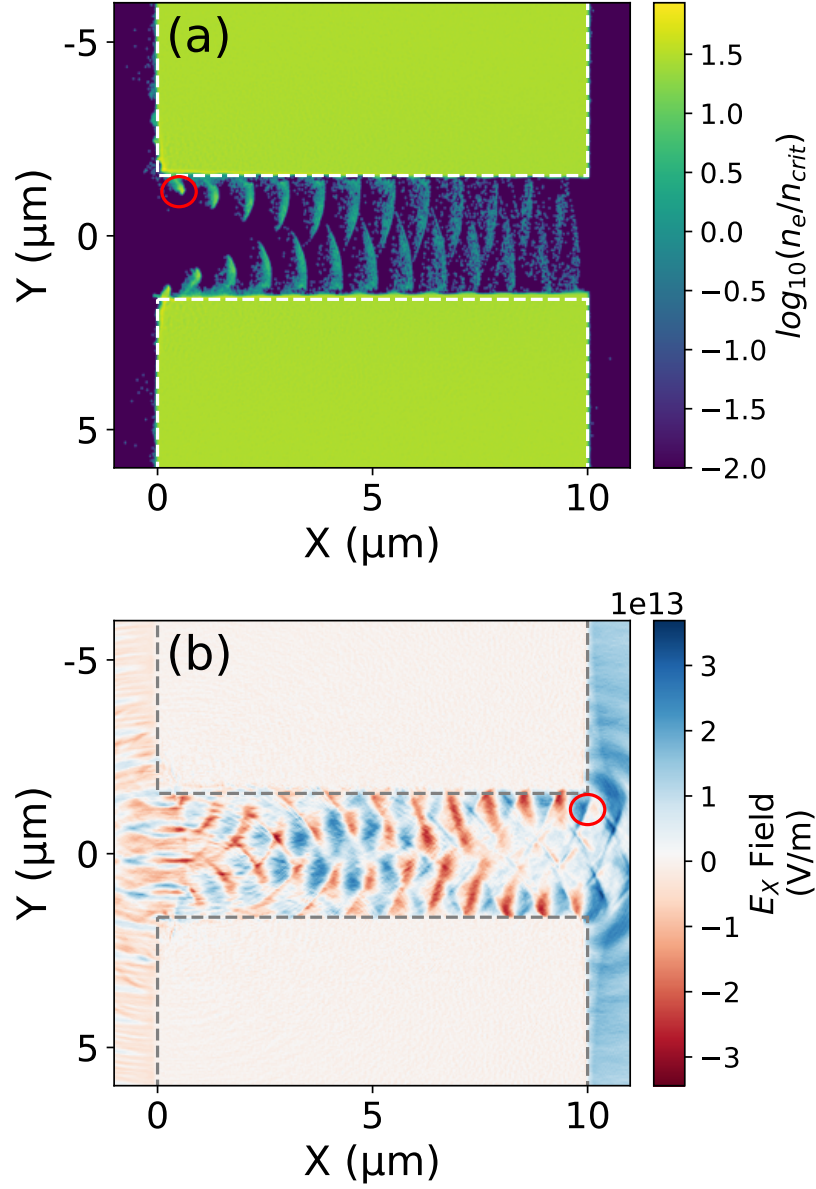


Figure 4.7: 2D simulation results for a target with $l = 10 \mu\text{m}$ and $d = 3 \mu\text{m}$: (a) electron density at $t = -4.8$ fs; (b) longitudinal electric field at $t = 2.8$ fs. The dashed lines indicate the initial positions of the target surfaces and aperture edges. An arbitrary electron bunch is circled in red for the purpose of tracking interaction metrics over time.

Since $2\omega_L$ light is generated from a single chain of electron bunches, any spatial structure observed in the emitted light can be attributed to an underlying structure in the electron bunch distribution as they are generated from the aperture edges. In the interactions shown in figures 4.5(a) and 4.7(a), the electron bunches exhibit a bipolar structure, with bunches being emitted from opposite edges of the aperture. This results in a spatial distribution of $2\omega_L$ light characterized by two distinct lobes, corresponding to

the positions on the aperture edges from which the electron bunches originate. The subject of spatial mode generation will be covered in more detail in later sections reporting on the use of 3D simulations.

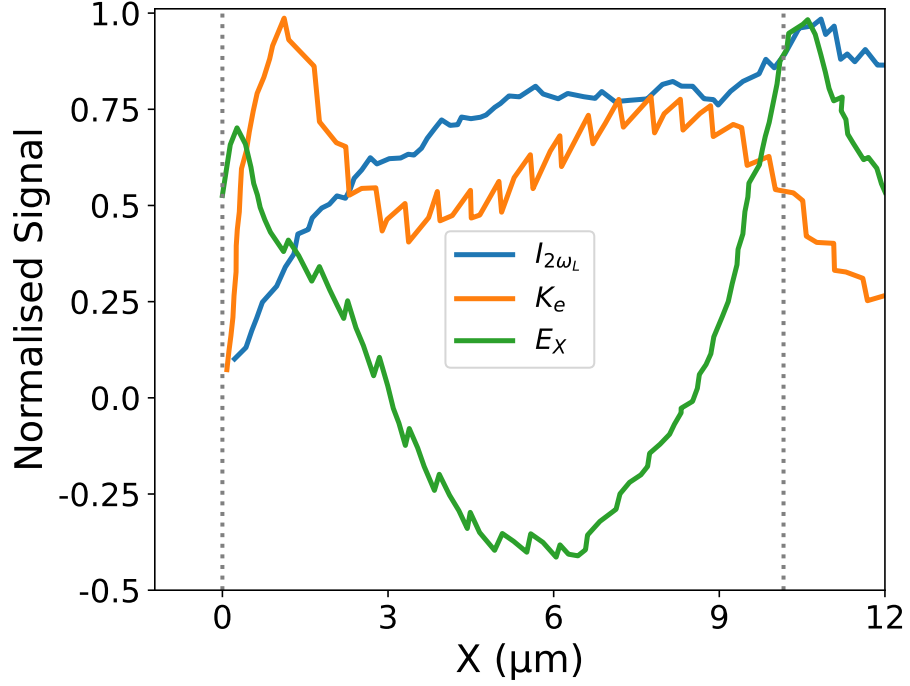


Figure 4.8: Total energy (K_e , orange) in the electron bunch marked by the red circle in figures 4.7(a) and (b) as it propagates in space, together with the intensity of the $2\omega_L$ light generated, averaged 400 nm around the bunch ($I_{2\omega_L}$, blue), and the longitudinal electric field experienced by the bunch (E_X , green) averaged over the same area. Note that arbitrary units are used for the quantities to allow for direct comparison.

The corresponding case for a target with $l = 10 \mu\text{m}$ and $d = 3 \mu\text{m}$ is shown in figures 4.7(a) and (b). Under these conditions, the aperture more closely resembles a channel. Figure 4.7(b) shows the electric field strength at a simulation time of $t = 24 \text{ fs}$, selected as this is the moment at which the electron bunch circled in figure 4.7(a) reaches the rear of the aperture channel. Examining the acceleration of this bunch reveals behaviour initially similar to the $l = 0.5 \mu\text{m}$ case, as shown in figure 4.8. However, within the aperture channel, the continued co-propagation of the electron bunch and laser pulse results in deceleration due to interactions with the positive longitudinal fields of the focused Gaussian pulse (arising from the divergence of the laser's electric field profile, i.e. $\nabla \cdot \vec{E} = 0$ in free space). Despite this deceleration, the electron bunch maintains relativistic energies throughout its propagation within the aperture channel. As the bunch continues to propagate, it de-phases from the laser pulse and encounters

a repeating positive and negative longitudinal field structure, visible in figure 4.7(b). This results in alternating acceleration and deceleration, as seen in figure 4.8 for $X > 3 \mu\text{m}$, until the bunch reaches the end of the aperture channel. Electrons within the bunch emit during both the acceleration and deceleration phases, with interference between this and emittance from the surrounding electron bunches generating a strong $2\omega_L$ signal. After exiting the aperture channel at $X > 10 \mu\text{m}$, the electron bunch interacts with the sheath field at the rear of the target, decelerating and thus leading to further $2\omega_L$ generation.

4.2.3 Investigating the interactions conditions for optimal $2\omega_L$ generation

Having investigated the physics underpinning the generation of frequency-doubled light, the next consideration was the effect of varying the target parameters of aperture diameter and target thickness on its production. A systematic parameter scan was performed using 2D EPOCH simulations during which l was varied from 0.5 to 10 μm and d was varied from 0.5 to 5 μm . The aim of this parameter scan was to determine the target thickness and aperture diameter that would optimise the conversion efficiency of drive pulse energy to frequency-doubled light. This efficiency was calculated from the ratio of the total energy of the generated $2\omega_L$ light detected 5 μm from the rear of the target to the total energy present in the input laser pulse. This process was simplified through the use of the BISHOP code [100] that automated the incremental changing of the parameter values within the simulations.

Figure 4.9 shows the conversion efficiency as it varies with aperture diameter and target thickness for a peak pulse intensity of $I_L = 1 \times 10^{21} \text{ W/cm}^2$. Labelled as region I are the conditions for optimal efficiency, which occur at an aperture diameter of $d \sim 2 \mu\text{m}$ (for $\phi_L = 2.1 \mu\text{m}$). For this aperture diameter, the efficiency saturates for a target thickness $\geq 6 \mu\text{m}$. Further increasing the aperture diameter results in less interaction between the aperture edges and the longitudinal field of the laser pulse, reducing the conversion efficiency, labelled as region II.

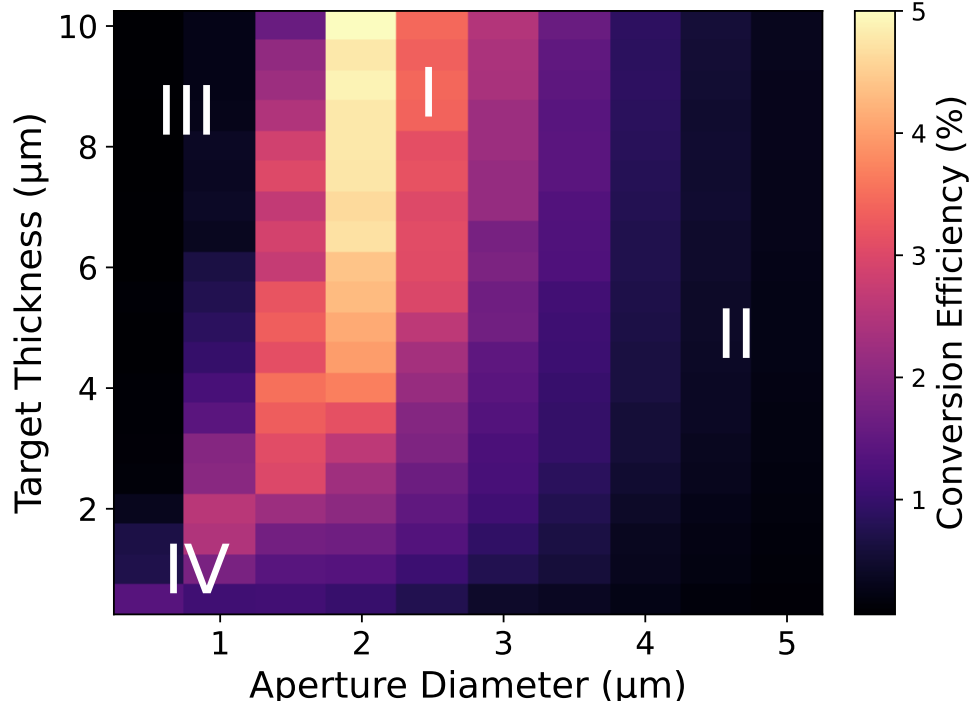


Figure 4.9: 2D simulation results showing the conversion efficiency to $2\omega_L$ light generation as a function of target thickness and aperture diameter, after the interaction of a linearly polarised laser pulse with an intensity of $I_L = 1 \times 10^{21} \text{ W/cm}^2$.

As the aperture diameter is reduced below the optimal $2 \mu\text{m}$, a decrease in conversion efficiency is observed for increasing target thickness, labelled as region III. This can be attributed to the poor matching conditions between the free-space Gaussian laser beam focus and the near cut-off modes that can exist within the aperture when its thickness exceeds the laser wavelength. This leads to the attenuation of the beam as it propagates through the aperture. For aperture diameters $< 2 \mu\text{m}$, the efficiency is found to be dependent on the target thickness. For thicker targets, the beam attenuation that occurs leads to a reduction in the efficiency to negligible levels. However, as labelled in region IV, appreciable efficiency can be recovered if the target thickness is reduced. For sufficiently thin targets, ($< 2 \mu\text{m}$) the transverse ponderomotive force that is generated by the focused light is great enough to expel the plasma from the focal-spot region to the degree that the poor matching conditions present for small apertures with greater thicknesses no longer exist. The result is an effective increase in the aperture diameter, allowing for the continued propagation of the pulse and generation of $2\omega_L$ light, similar to that of the interaction with an ultra-thin target as reported by Duff *et al* [56].

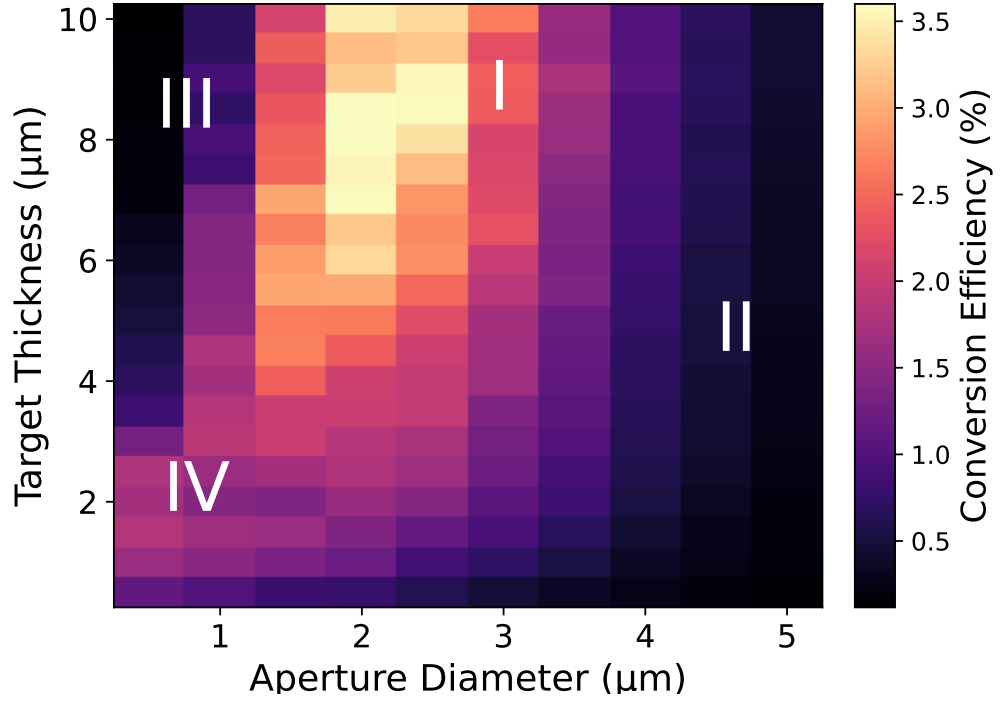


Figure 4.10: 2D simulation results showing the conversion efficiency to $2\omega_L$ light generation as a function of target thickness and aperture diameter, after the interaction of a linearly polarised laser pulse with an intensity of $I_L = 1 \times 10^{22} \text{ W/cm}^2$.

A similar parameter scan in which l and d were varied over the same range of values was then performed in which the peak pulse intensity increased to $I_L = 1 \times 10^{22} \text{ W/cm}^2$, with the results displayed in figure 4.10. The effect of plasma expulsion from the aperture becomes more pronounced at this higher intensity, evidenced by the larger range of target thicknesses over which generation efficiency is observed for the smallest aperture diameters, reflected in the widening of region IV. This is also likely the cause of the slight increase of the range of aperture diameters covered by region I, indicating that the aperture diameter for optimal efficiency now includes $\sim 2.5 \mu\text{m}$.

The final parameter scan increased the peak pulse intensity to $I_L = 1 \times 10^{23} \text{ W/cm}^2$, with the results shown in figure 4.11. The resulting increase in transverse ponderomotive force that accompanies the increased pulse intensity leads to the broadening of region IV across the entire range of target thicknesses for the smallest aperture diameters, resulting in the absence of region III. The reduced peak density used for the simulations is likely to have exaggerated this effect.

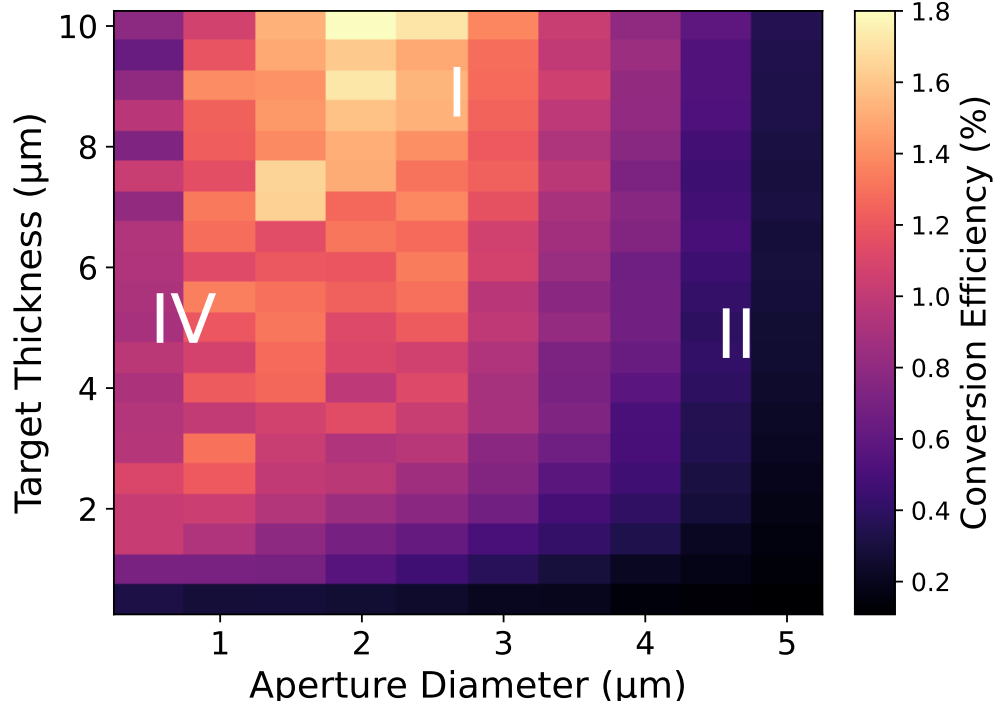


Figure 4.11: 2D simulation results showing the conversion efficiency to $2\omega_L$ light generation as a function of target thickness and aperture diameter, after the interaction of a linearly polarised laser pulse with an intensity of $I_L = 1 \times 10^{23} \text{ W/cm}^2$

Under these conditions, the edges of the apertures will undergo sufficient heating during the interaction to become relativistically transparent as the pulse intensity approaches its maximum, independent of the initial aperture diameter. The subsequent increase in effective aperture diameter relative to its initial size acts to reduce the efficiency of the generation mechanism. This effect is responsible for the reduction in maximum conversion efficiency observed with increasing pulse intensity, as the effective aperture diameter is larger at such intensities, departing further from the optical size. The aperture diameter under which the optimal efficiency is observed is the same as that of the case of the $I_L = 1 \times 10^{22} \text{ W/cm}^2$.

4.3 Efficiency modelling

In an attempt to better understand the dependence of $2\omega_L$ light generation on the parameters of target thickness, a model of the conversion efficiency scaling with these parameters was constructed. The generation mechanism put forward in the preceding sections posits that the electron bunches are accelerated and decelerated by the lon-

itudinal electric field of the pulse as it interacts with the aperture edges. It is then assumed that the magnitude of this field near the aperture edge plays a role in determining the mechanism's conversion efficiency of ω_L to $2\omega_L$ light. For a pulse with a focused Gaussian spatial profile, the longitudinal electric field can be found by taking the spatial differential of the electric field along the polarisation direction. This field, E_X , is defined as;

$$E_X(Y) = \frac{d}{dy} e^{-\left(\frac{2 \cdot \sqrt{\ln(2)} \cdot Y}{\phi_L}\right)^2} \quad (4.2)$$

for a pulse that is linearly polarised in the Y -axis. The transverse position at which this field reaches its maximum absolute value is at $Y \sim \phi_L/2$. As such, with increasing aperture diameter, the magnitude of this longitudinal electric field at the aperture edge, E_X , is maximised when $d \sim \phi_L$; this is shown in figure 4.12(a). Note that the field is modelled simply as the derivative of the Gaussian spatial profile as the purpose of the model is to investigate trends in the conversion efficiency rather than predict absolute magnitudes. A full treatment of the electric fields associated with a Gaussian spatial pulse is given in [101].

Another key factor influencing $2\omega_L$ light generation is the percentage of the laser pulse that can propagate into the aperture. Achieving higher efficiency requires a larger proportion of the available light to successfully enter and pass through the aperture. For the 2D simulations, this quantity, K_{in} , is described by the following equation;

$$K_{in}(d) = \int_{-d/2}^{d/2} e^{-\left(\frac{2 \cdot \sqrt{\ln(2)} \cdot Y}{\phi_L}\right)^2} dY \quad (4.3)$$

For smaller aperture diameters, a greater proportion of the pulse is blocked by the bulk of the target, hence this equations dependence on d . This is shown in figure 4.12(b).

As this is a relatively simple model, the role that relativistic transparency plays in increasing the aperture diameter, and so the percentage of the light that can enter the aperture, is not included as this would be considered a second-order effect.

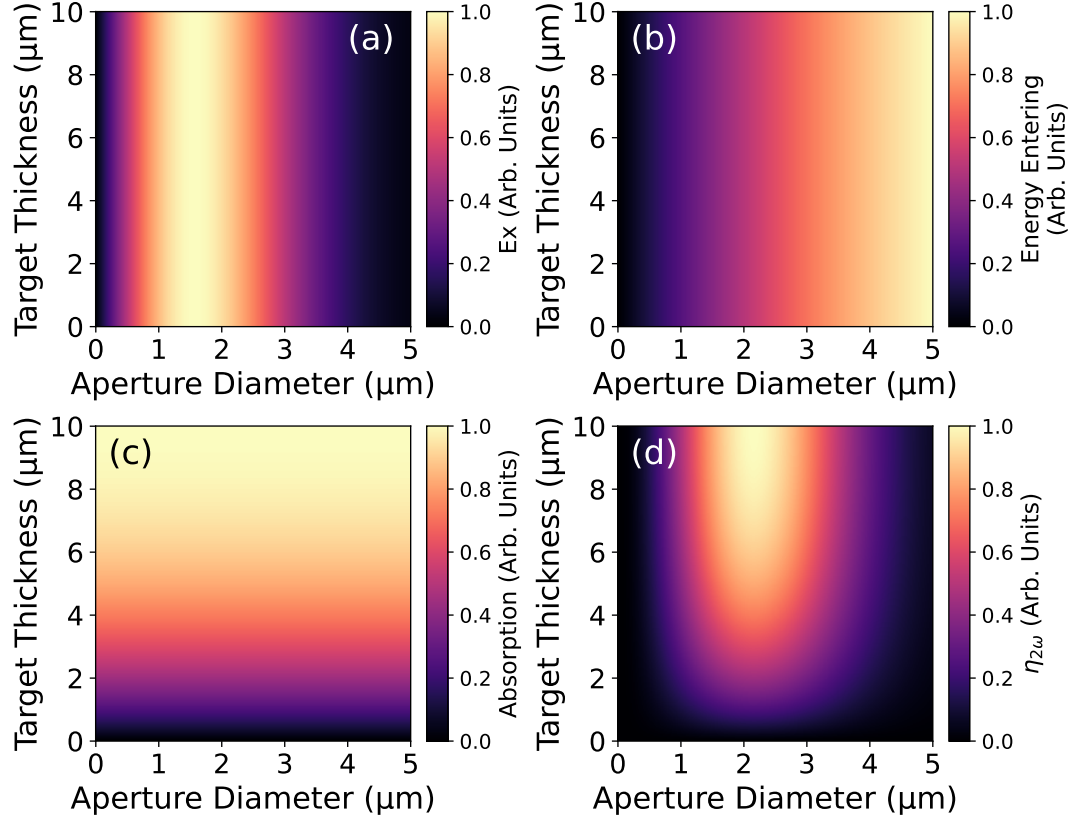


Figure 4.12: Three parameters that vary with both target thickness and aperture diameter: (a) magnitude of the longitudinal electric field at the aperture edge; (b) the energy into the aperture; (c) the laser absorption. (d) conversion efficiency to $2\omega_L$, as defined by equation . 4.5.

The final parameter considered for this model is that of the absorption of the pulse energy into the electron bunches that are accelerated from the aperture edges, here labelled Abs_L . With the assumption that the number of electrons exposed to the laser pulse increases linearly with target thickness, an absorption coefficient α is defined. A second assumption is that the laser pulse decays exponentially as it propagates through the aperture, with an attenuation coefficient defined as β . Laser pulse attenuation is likely affected by various factors, including self-focusing and reflection. Therefore, the attenuation constant is treated as a free parameter and is approximated based on the levels of transmitted light observed in the simulation results. The absorption of the pulse is then given by the following expression;

$$Abs_L(l) = \alpha l e^{-\beta l} \quad (4.4)$$

where l is target thickness. This is shown in figure 4.12(c).

Combining the three factors considered for this model gives the following first order

approximation for the $2\omega_L$ conversion efficiency;

$$\eta_{2\omega} = |E_X(d/2)|\text{Abs}_L(l)K_{in}(d) \quad (4.5)$$

Note that this is based on the assumption that the parameters in this expression are related linearly. Plotting this relationship with d and l , shown in figure 4.12(d), provides a good agreement with the simulation results shown in figure 4.9. This simple model serves mainly to illustrate the general trend in conversion efficiency as it varies with target thickness and aperture diameter. It does not account for factors such as relativistic transparency and reabsorption of generated $2\omega_L$ light. The former primarily affects only the thinnest targets and highest-intensity pulses, while the latter becomes significant only for very thick targets.

4.4 3D Simulations

The next stage in the investigation involved the use of 3D EPOCH simulations, allowing for the scaling of conversion efficiency with the parameters l and d to be tested under more realistic conditions. This also enabled the investigation of the change in mode shape with pulse polarisation, as harmonic light has been shown to be produced in non-symmetric modes [56][62] that are not fully observable in 2D simulations. For these simulations, the XYZ grid was defined as $1000 \times 720 \times 720$ cells, corresponding to $20 \mu\text{m} \times 20 \mu\text{m} \times 20 \mu\text{m}$. The targets were, as with the 2D simulations, comprised of a slab of Al^{13+} ions, neutralised with an electron population with a peak density of $30n_{crit}$ and an initial temperature of 10 keV. All other simulation parameters were kept the same as in the 2D case, aside from the particle-per-cell number being reduced to 5 in an effort to reduce the computational resources required to perform these simulations. Despite this, the cost of running the 3D simulations remained significantly higher than in the 2D case, leading to a courser parameter scan being performed to confirm whether the efficiency trends identified in 2D simulations were also observed in 3D.

Figure 4.13 shows the results of the 3D parameter scan. The conversion efficiencies dependence on the target parameters l and d are shown to be consistent with those observed with the 2D simulations. Notable differences, however, are the minor increase in the aperture diameter at which the maximum efficiency is achieved, $(3 \pm 0.5) \mu\text{m}$ (compared with $2 \mu\text{m}$ for 2D), and the magnitude of the optimal efficiency reducing to

3.5%. The latter can be attributed to the fact that in 3D there is the possibility of the polarisation vector being parallel to the aperture edge, in which case there is minimal interaction. This is not possible in the case of a 2D simulation as the laser intensity only exists in the polarisation plane, which is perfectly normal to the aperture edges. This then leads to a portion of the pulse energy not being available to the generate $2\omega_L$ light in the case of a 3D simulation. Furthermore, it may be that the added degree of freedom present in the 3D simulations allows electrons to no longer be strictly confined to the laser fields when undergoing acceleration, reducing the overall number of electrons that are contributing to $2\omega_L$ generation.

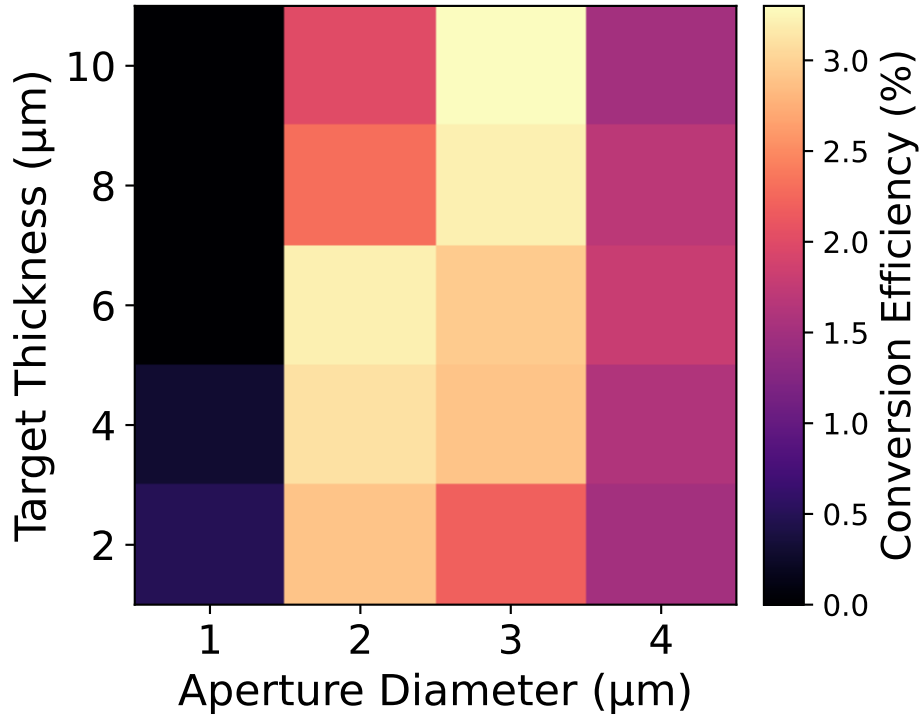


Figure 4.13: 3D simulation results showing the conversion efficiency to $2\omega_L$ light generation as a function of target thickness and aperture diameter, after the interaction of a linearly polarized laser pulse with an intensity of $I_L = 1 \times 10^{21}$ W/cm².

Returning to the efficiency modelling, for the expression defining the total energy entering the aperture (4.3) to be valid in 3D, two new factors must be included. The first is area of the aperture, rather than just a linear distance, and the second is the total area of the laser focus. The revised expression for $K_{in}(d)$ is as follows

$$K_{in}(d) = \int_{-\pi}^{\pi} \int_0^{d/2} e^{-\left(\frac{2\sqrt{\ln(2)} \cdot r}{\phi_L}\right)^2} r dr d\phi \quad (4.6)$$

where the integral is given in cylindrical coordinates.

Figure 4.14 shows the conversions efficiency, $\eta_{2\omega}$, as a function of l and d , using equation 4.6 to account for the changes in 3D. The updated efficiency modelling shows an increase in the optimal d by $\sim 0.5 \mu\text{m}$, agreeing with the 3D simulations.

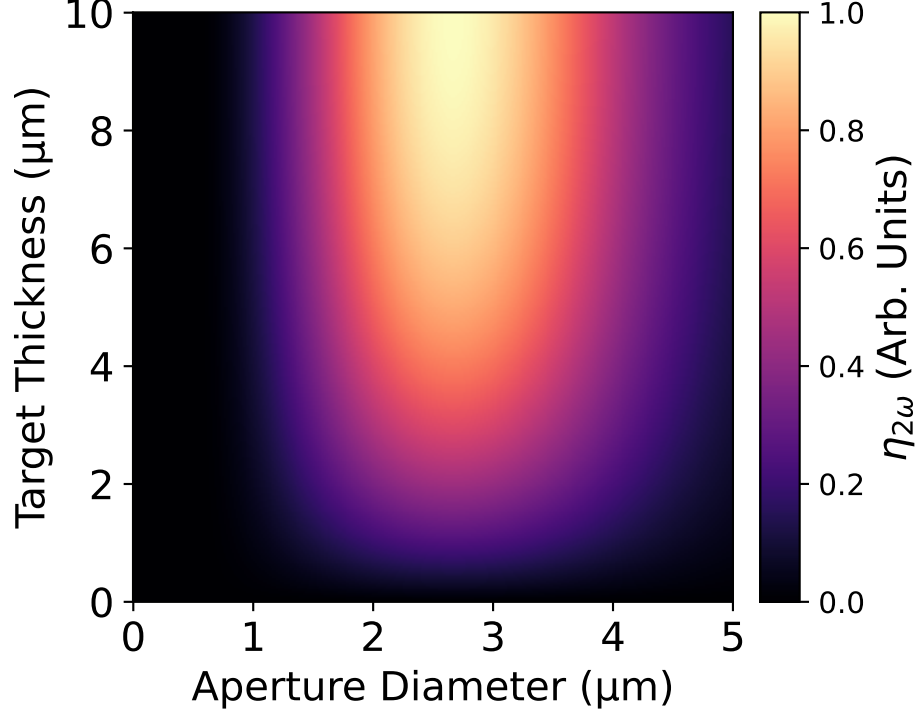


Figure 4.14: Conversion efficiency to $2\omega_L$ under 3D conditions for a pulse with $I_L = 1 \times 10^{21} \text{ W/cm}^2$ as it varies with both target thickness and aperture diameter, as defined by equation 4.5, amended to include equation 4.6.

Figure 4.15 compares the results of 2D and 3D simulations for the conversion efficiency dependencies on (a) aperture diameter with a fixed target thickness of $l = 6 \mu\text{m}$, and (b) target thickness at the optimal aperture diameters ($d = 2 \mu\text{m}$ for 2D and $d = 3 \mu\text{m}$ for 3D). The corresponding dependencies calculated from the efficiency modelling are also included. Overall, the results demonstrate good agreement between the simulations and the efficiency modelling. The simulations performed to investigate the case of circular polarisation are discussed later in the chapter.

With respect to the simulations, a rapid increase in efficiency is seen as d is increased to the optimum value, after which a slower decrease with aperture diameter is seen in the 3D case. As predicted by the efficiency modelling, a slight increase in the optimum aperture diameter is observed when comparing 2D and 3D simulations, rising from $2 \mu\text{m}$ to $3 \mu\text{m}$ respectively. The efficiency is also seen to increase at a slightly faster rate

in the 3D case when target thickness is varied, saturating at 3.5%. The slight difference in dependencies between the two simulation cases may be attributable to the optimal d being slightly larger in the 3D case.

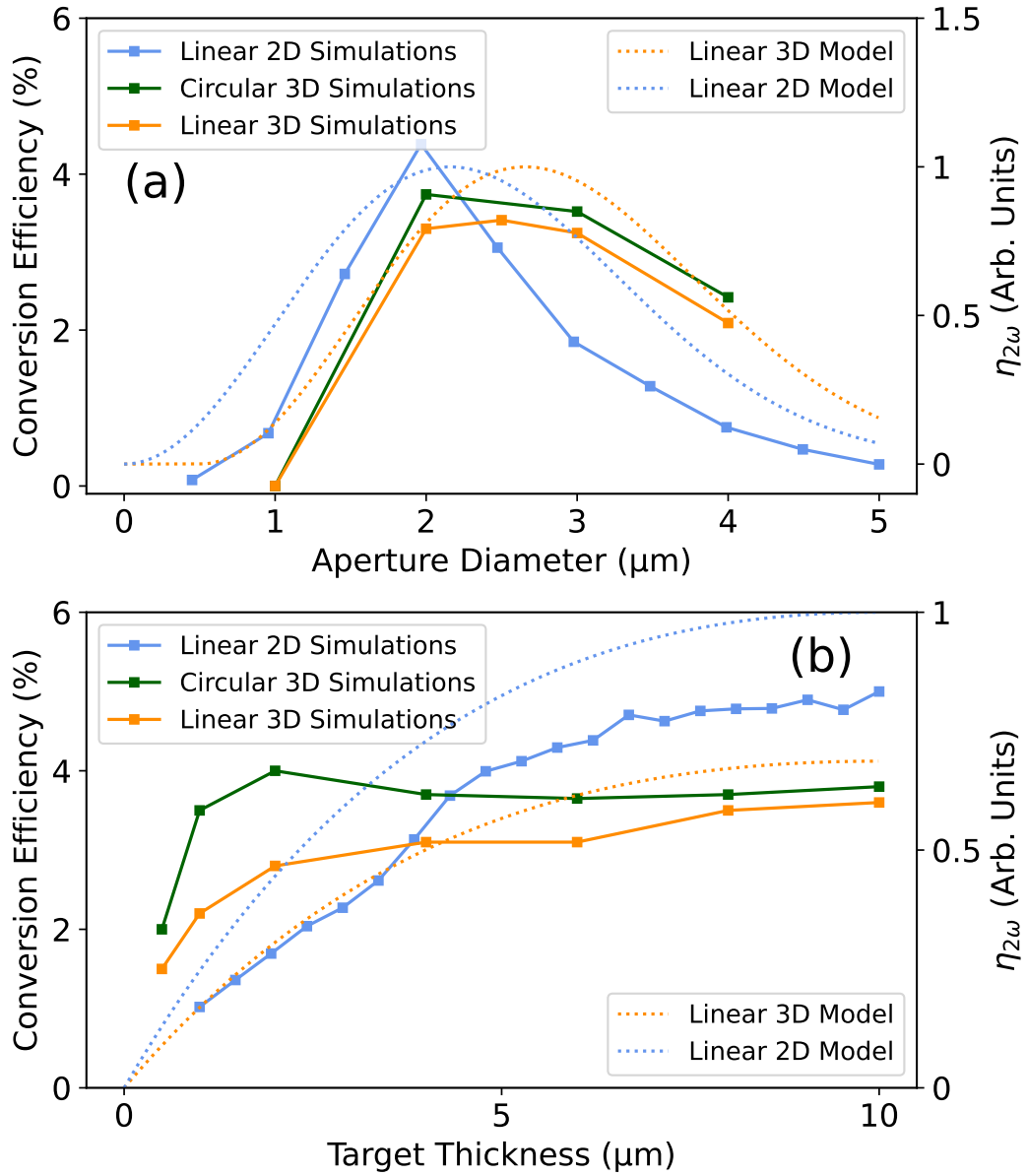


Figure 4.15: 3D simulation results showing the $2\omega_L$ generation efficiency as a function of (a) d for $l = 6 \mu\text{m}$, and (b) l for $d = 3 \mu\text{m}$. 3D linear polarisation results are shown in blue and circular polarisation results are shown in orange. 2D simulation data are included for comparison (dashed black line) for fixed $l = 6 \mu\text{m}$ and $d = 2 \mu\text{m}$ in (a) and (b), respectively. $I_L = 1 \times 10^{21} \text{ W/cm}^2$ for all cases.

Comparing the expansion of the laser pulse after passing through focus may also provide insight into dimension-related effects. In 2D, the intensity decreases with r , as opposed to with r^2 in 3D, acting to change the effective intensity. Finally, second-order effects that can change the intensity, and thus the longitudinal electric field, may be responsible

for difference between 2D and 3D. Examples of such effects include reflection of the laser pulse and self-focusing [102][103] as the pulse propagates through the aperture.

Figure 4.16 illustrates scalloping and self-focusing behaviour within the plasma aperture, with the drive laser pulse's intensity profile displayed for two targets with differing thickness. Taken from 3D simulations in which a target with $l = 0.5 \mu\text{m}$ was used, figures 4.16(a) and (b) show the ω_L intensity profile along $Z = 0 \mu\text{m}$ in the XY plane at simulation times of $t = 16$ and $t = 32$ fs respectively.

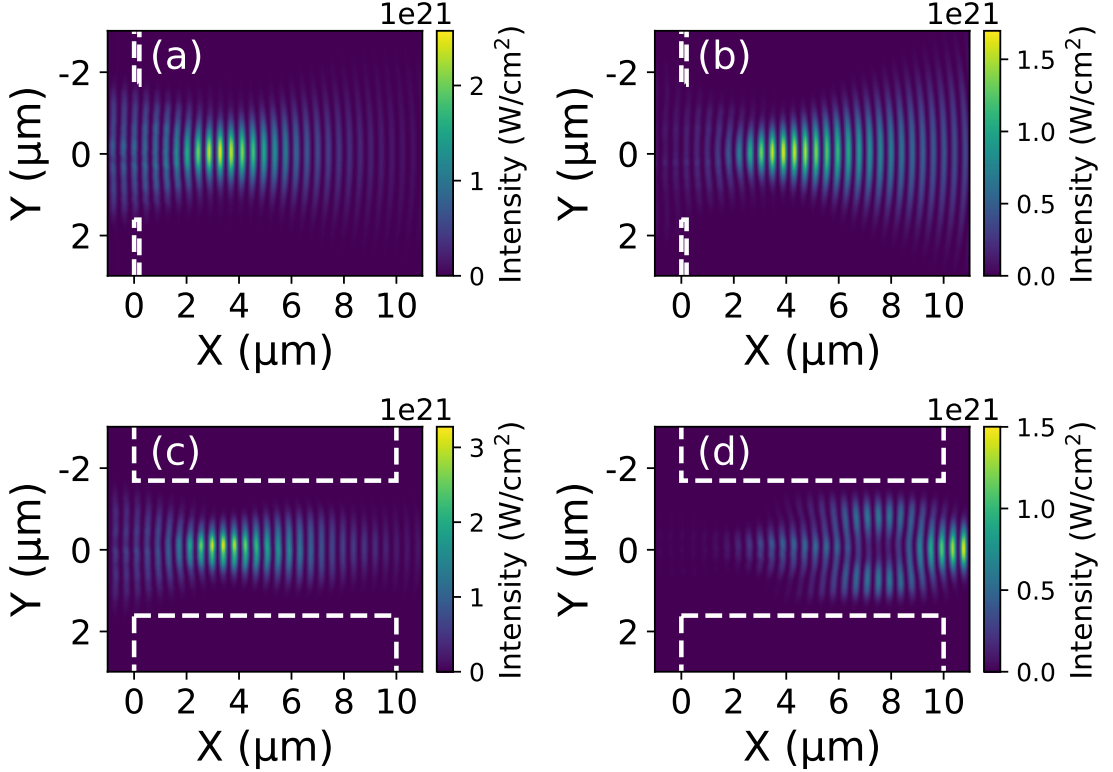


Figure 4.16: Intensity profile of the drive laser light in the XY plane at $Z = 0$ from 3D simulations, for $d = 3 \mu\text{m}$ and $l = 0.25 \mu\text{m}$, at: (a) $t = 16$ fs and (b) $t = 32$ fs. (c) and (d) are the corresponding results for $l = 10 \mu\text{m}$. The dashed white lines indicate the initial target profile and the nominal focused intensity of $I_L = 1 \times 10^{21} \text{ W}/\text{cm}^2$ in both cases.

As the pulse propagates through the aperture, it is seen to experience self-focusing to a position of $X \sim 3 \mu\text{m}$ before expanding into free space at the rear of the target. This self-focusing leads to an increase in the pulse intensity from the nominal value of $1 \times 10^{21} \text{ W}/\text{cm}^2$ (without plasma) to that of $2.5 \times 10^{21} \text{ W}/\text{cm}^2$, a factor of two greater than that achievable were the same focus used in 2D simulations. This can be attributed to dimensionality effects; the pulse energy can be focused radially in the case of 3D, but only in a plane in 2D. Shown in figures 4.16(c) and (d) is the ω_L intensity

profile at the same simulation times for a target of thickness $l = 10 \mu\text{m}$. Similar self-focusing behaviour is observed, although with a higher maximum intensity of $\sim 3 \times 10^{21} \text{ W/cm}^2$ achieved and an inhibiting of pulse expansion due to the increased length of the aperture channel. The pulse is seen to have refocused toward the rear of the target in the later time-step shown in figure 4.16(d).

4.4.1 Generated $2\omega_L$ spatial structure

An important aspect of the generated $2\omega_L$ light that has not yet been examined is its spatial structure. While the cause, linked to the electron bunch structure as they are accelerated from the aperture edges, was briefly discussed in section 4.2.2, a detailed investigation of the spatial structure requires the use of 3D simulations. This section explores the dependence of the generated spatial modes on interaction parameters, specifically pulse polarisation and aperture diameter, through the use of such simulations. The goal is to enable tailoring of the mode characteristics for their use in various applications, such as optimising focusing behaviour in laser-plasma interactions to enhance particle acceleration [104], and generating radiation sources with orbital angular momentum (OAM)[105].

Pulse polarisation

Variability in the $2\omega_L$ spatial structure that can be generated was first investigated through changing the drive pulse polarisation. Figure 4.17(a) and (b) show the $2\omega_L$ time-integrated spatial mode measured $5 \mu\text{m}$ from the rear of the target for linearly and circularly polarised light respectively. Target parameters of $d = 4 \mu\text{m}$ and $l = 6 \mu\text{m}$ and a pulse intensity of $I_L = 1 \times 10^{21} \text{ W/cm}^2$ were used.

The radiant fluence is plotted on a log scale in these figures to demonstrate that there are also low amounts of $2\omega_L$ light produced outside of the central mode in both polarisation cases, with even higher-order structures. This is caused by the interference of the generated light varying when the sources are sufficiently far apart [62].

For a linearly polarised drive pulse, a TEM_{01} mode is observed, while a circularly polarised pulse results in an annular intensity profile. An important feature of the light generated from the use of a circularly polarised pulse is the conversion of the spin angular moment of the ω_L light into $2\omega_L$ light possessing OAM [62][106].

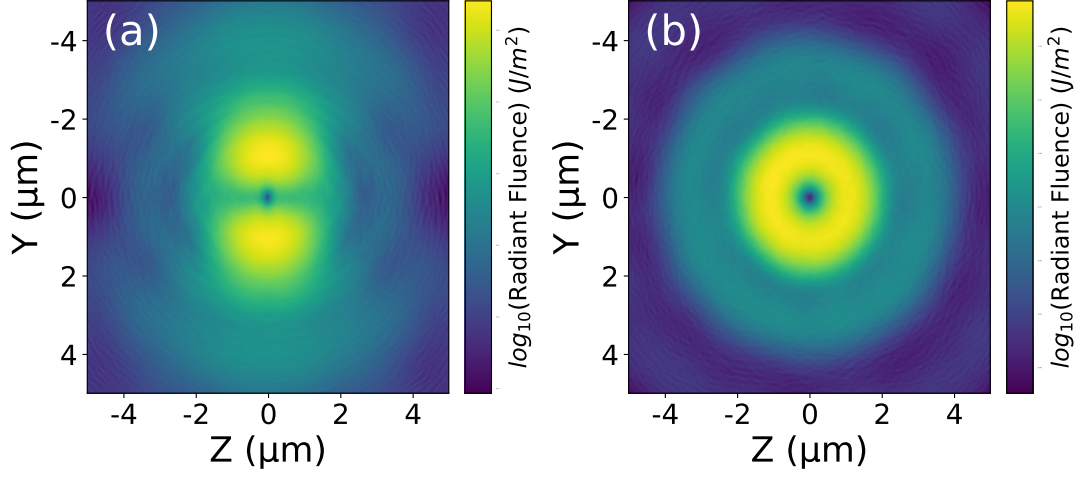


Figure 4.17: $2\omega_L$ time-integrated spatial mode structure at $5\ \mu\text{m}$ from the rear of the target with $d = 4\ \mu\text{m}$ and $l = 6\ \mu\text{m}$, for (a) linear polarisation and (b) circular polarisation. Fluence is plotted on a logarithmic scale to show the higher-order spatial structures present at larger aperture radii.

A key factor in understanding the influence of pulse polarisation on the spatial structure is the difference between the time-integrated profile and that observable on a sub-laser-cycle time scale. For a linearly polarised pulse, the spatial structure of the $2\omega_L$ light on a sub-laser-cycle timescale would reveal a structure similar to that shown in figure 4.17(a), as the positions on the aperture edges from which electron bunches are accelerated remains stable, assuming minimal change in aperture geometry throughout the interaction. In contrast, with circularly polarised light, the spatial structure exhibits rotational behaviour on sub-laser-cycle timescales, resulting in a helical structure over time.

This behaviour can be illustrated by examining the simulation from which the time-integrated mode shown in figure 4.17(b) was produced. Shown in figure 4.18(a) is the electric field strength measured $5\ \mu\text{m}$ from the rear of the target at a simulation time of $t = 540\ \text{fs}$. Then shown in figures 4.18(b), (c) and (d) are the electric field strengths for simulation times of $t + \frac{T}{4}$, $t + \frac{T}{2}$ and $t + \frac{3T}{4}$ respectively, with T being the laser period. As can be seen, there is a full rotation of the spatial structure of the light over this duration. This variation in mode shape arises from the rotation of the polarisation vector that occurs with circularly polarised light as it propagates. This rotation results in the polarisation vector remaining perpendicular to the aperture edge, causing the acceleration of electron bunches from the aperture edges in a helical pattern. When time-integrated, this helical structure produces the annular profile observed in figure 4.17(b).

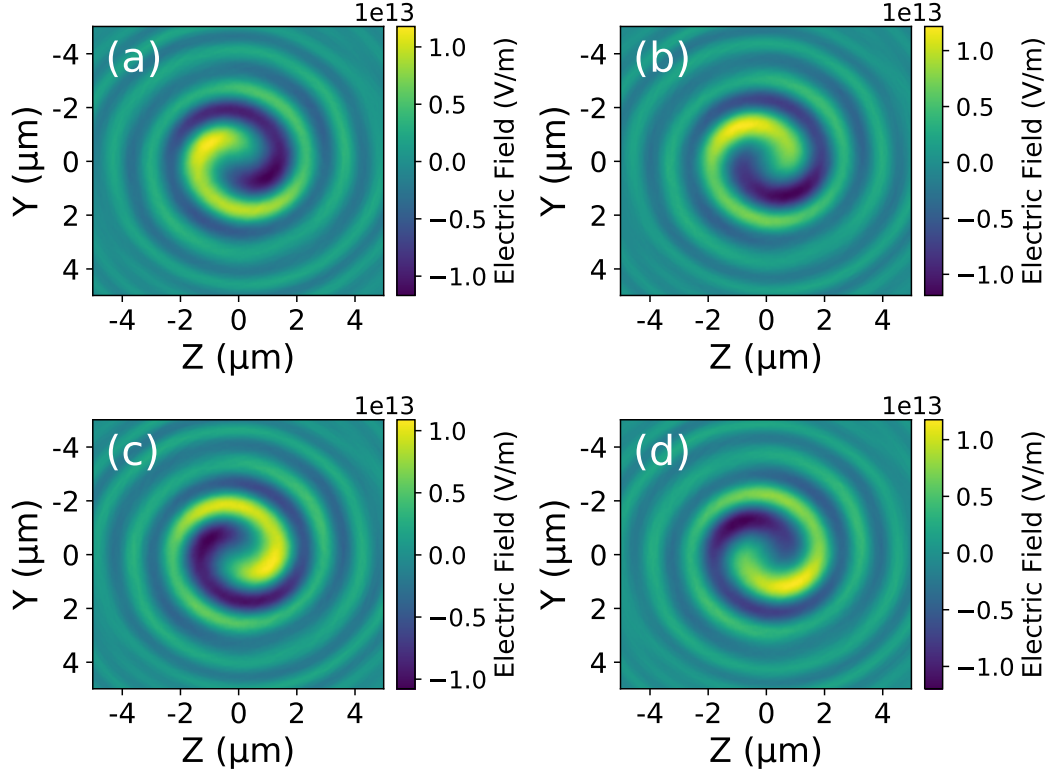


Figure 4.18: Electric field strength measured $5 \mu\text{m}$ from the rear of a target with $d = 4 \mu\text{m}$ and $l = 6 \mu\text{m}$ for a circularly polarised pulse at simulation times: (a) $t = 540 \text{ fs}$, (b) $t + \frac{T}{4}$ (c) $t + \frac{T}{2}$ and (d) $t + \frac{3T}{4}$. Here, T represents the period of the drive laser pulse, with an incident pulse intensity of $I_L = 1 \times 10^{21} \text{ W/cm}^2$.

Aperture diameter

In addition to its dependence on polarisation, the change in mode shape with aperture diameter was investigated. Two sets of 3D simulations were conducted, one with linearly polarised light and the other with circularly polarised light, in which the aperture diameter was varied from 1 to $4 \mu\text{m}$. Target thickness was fixed at $6 \mu\text{m}$ for both sets of simulations.

The change in mode size with aperture diameter for linearly polarised light is displayed in figure 4.19, with the mode size defined as the distance between central points of the two lobes of generated $2\omega_L$ light. Figures 4.20(a) through (d) show the time integrated modes generated from the use of a pulse that is linearly polarised in the Y -axis for the aperture diameters used in the simulations.

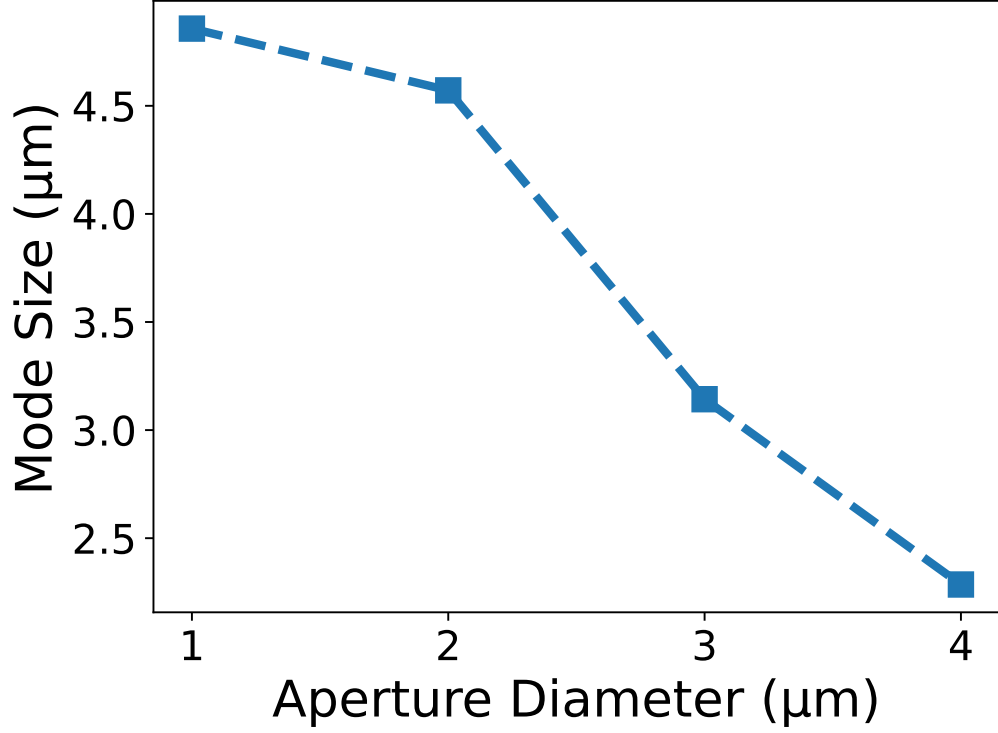


Figure 4.19: Mode size of the generated $2\omega_L$ light, measured $5 \mu\text{m}$ from the rear of the target, as a function of aperture diameter, for an $l = 6 \mu\text{m}$ target and a linearly polarised pulse with $I_L = 1 \times 10^{21} \text{ W/cm}^2$.

An inverse relationship is observed between the mode size and the aperture diameter, highlighting the potential to tailor the spatial profile of the generated $2\omega_L$ light. The spatial profiles reveal a clear variation in the levels of $2\omega_L$ light generated as the aperture diameter changes, with the lowest levels produced from the smallest aperture diameter ($d = 1 \mu\text{m}$). As the aperture diameter increases, the generated mode becomes more distinct, accompanied by a noticeable increase in the quantity of $2\omega_L$ light produced.

Figures 4.21 and 4.22 illustrate the variation in mode diameter and the time-integrated spatial profiles, respectively, for the simulations performed with circularly polarised light. Similar to the case of linear polarisation, an inverse relationship between the mode diameter and aperture diameter is observed. Additionally, as the aperture diameter increases, the quantity of $2\omega_L$ light generated also increases, with the mode becoming more distinct in the process.

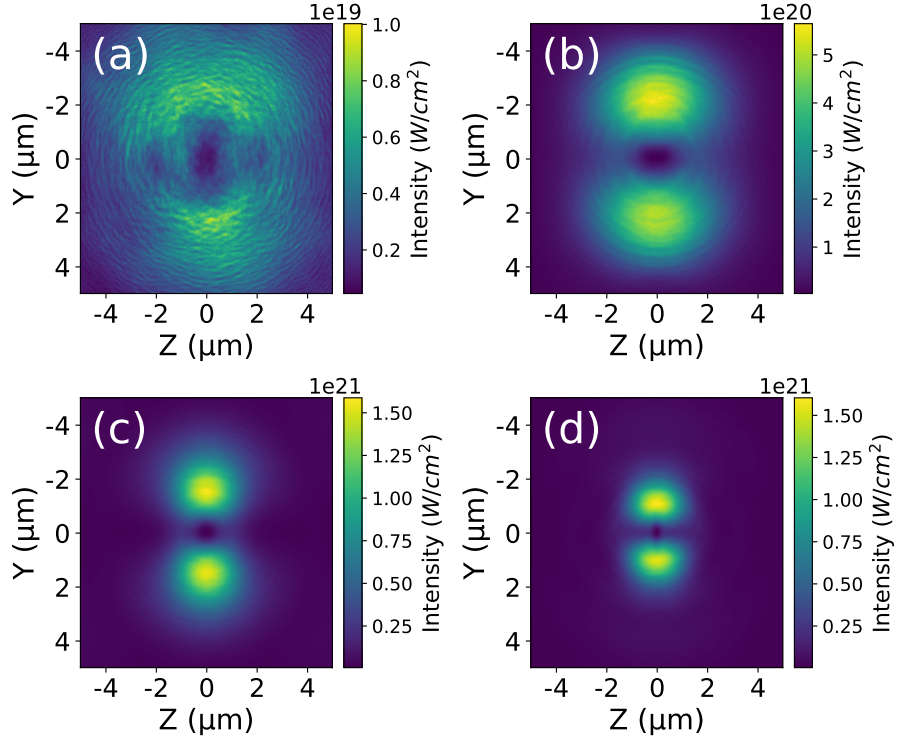


Figure 4.20: Time-integrated spatial mode structure at $5 \mu\text{m}$ from the rear of the target resulting from the use of a drive pulse linearly polarised in the Y -axis. Aperture diameters of (a) $1 \mu\text{m}$, (b) $2 \mu\text{m}$, (c) $3 \mu\text{m}$, and (d) $4 \mu\text{m}$, all for $l = 6 \mu\text{m}$ and $I_L = 1 \times 10^{21} \text{ W/cm}^2$.

Finally, when investigating the conversion efficiency as it varies with d and l for circular polarisation, the overall trends are similar to the case of linear polarisation, as shown in figures 4.15(a) and (b).

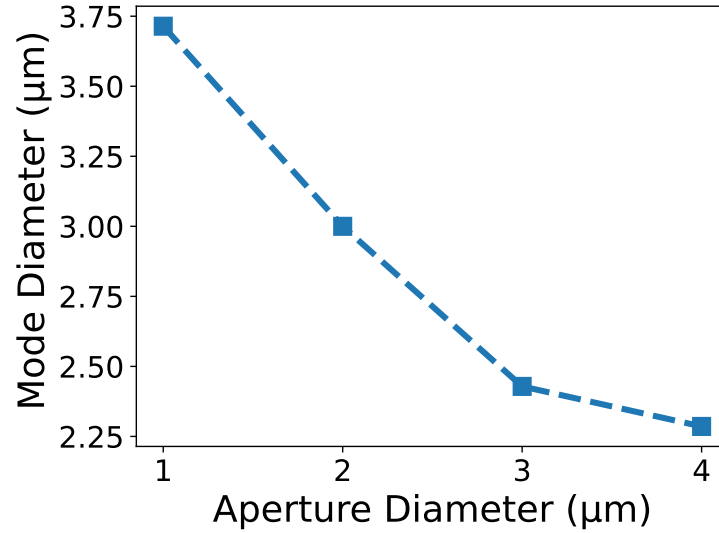


Figure 4.21: Mode diameter of the generated $2\omega_L$ light, measured $5 \mu\text{m}$ from the rear of the target, as a function of aperture diameter, for an $l = 6 \mu\text{m}$ target and a circularly polarised pulse with $I_L = 0.5 \times 10^{21} \text{ W/cm}^2$.

However, a difference is observed in the absolute conversion values, with this being attributed to the polarisation vector always being normal to the aperture edge, as covered previously in this section. This results in the constant generation of $2\omega_L$ light over the wave period.

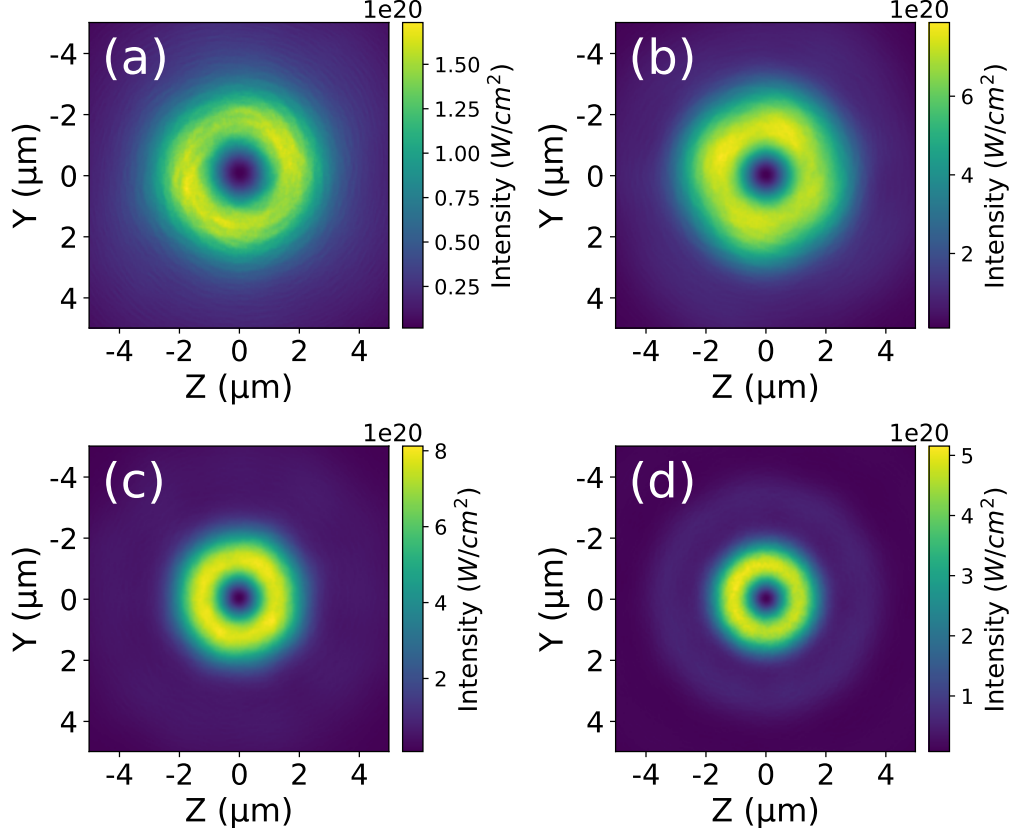


Figure 4.22: Time-integrated spatial mode structure at $5\ \mu\text{m}$ from the rear of the target resulting from the use of a circularly polarised drive pulse. Aperture diameters of (a) $1\ \mu\text{m}$, (b) $2\ \mu\text{m}$, (c) $3\ \mu\text{m}$, and (d) $4\ \mu\text{m}$, all for $l = 6\ \mu\text{m}$ and $I_L = 0.5 \times 10^{21}\ \text{W}/\text{cm}^2$.

The variation in mode size with aperture diameter for both polarisation cases arises from the corresponding change in the pulse's angular spread after propagating through the aperture channel. As the aperture diameter decreases, the angular spread increases, leading to a similar increase in the emission angle of electron bunches as they are accelerated.

4.5 Effects of temporal intensity contrast

A factor that has so far not been considered is the effect of the temporal intensity contrast of the laser pulse (3.2.2) on the generation of $2\omega_L$ light. A more comprehensive

investigation of the form that the pre-plasma in the aperture channel is reported on in section 5.2.1 but is beyond the scope of this chapter. As a simple approximation that would allow for the effects of pre-plasma to be studied, the choice was made to initialise 3D simulations with plasma of various densities filling the aperture channel, ranging from $0.01 \times$ to $1 \times n_{crit}$. While a complete and uniform filling of the aperture is not the exact form that the pre-plasma would take, the combination of the pre-pulse and rising edge of the main pulse would likely form an under-dense plasma in the aperture channel, and this profile is easily implemented within EPOCH. These simulations involved a target with a thickness of $6 \mu\text{m}$ and aperture diameter of $3 \mu\text{m}$.

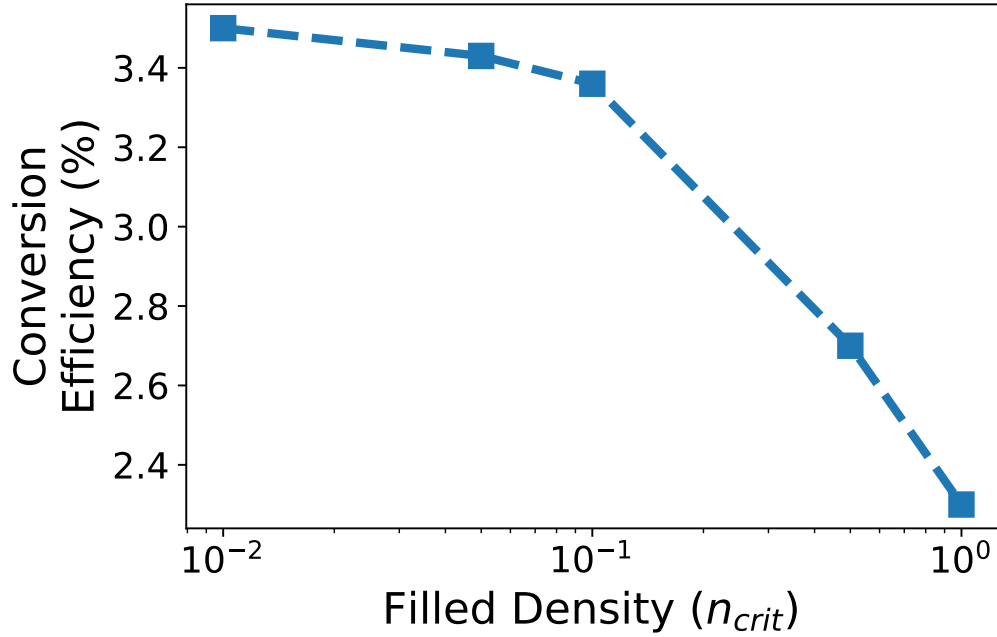


Figure 4.23: Conversion efficiency to $2\omega_L$ light as a function of the plasma density within the aperture, for a target with $d = 2 \mu\text{m}$, $l = 6 \mu\text{m}$ and pulse with $I_L = 1 \times 10^{21} \text{ W/cm}^2$.

Figure 4.23 shows the results of varying the plasma densities in the aperture channel on the conversion efficiency to frequency-doubled light. A notable reduction is observed at higher densities, with the efficiency decreasing to around 60% of the maximum value at $n_e = n_{crit}$. At lower densities, less than $0.1n_{crit}$, there is a reduction of only several percent observed, suggesting minimal impact on the generation of the $2\omega_L$ light. A possible cause of the reduction in conversion efficiency with increasing plasma density within the aperture channel is the hindering of the main pulse. As the pulse is prevented from propagating through the aperture in the initial stages of the interaction, the generation of frequency-doubled light may be suppressed. Once the

plasma has been forced from the aperture channel, $2\omega_L$ generation can occur. As the plasma density approaches n_{crit} , this ‘clearing’ of the channel could take longer, hence the inverse relationship between density and conversion efficiency. This relationship is discussed further in section 5.2.4.

4.6 Aperture tapering

All previous investigations of $2\omega_L$ light generation from aperture targets have involved geometries for which the entrance and exit diameters are equal in size, resulting in a cylindrical channel when the target’s thickness-to-diameter ratio is large. To explore the possibility of increasing conversion efficiency with alternative geometries, targets for which the entrance and exit diameters were not equal, creating a tapered aperture channel, were considered. This was investigated via a set of 3D simulations in which this target geometry was achieved by defining two aperture diameters that represented the entrance and exit diameters, labelled d_1 and d_2 respectively. Keeping d_2 fixed at $2\text{ }\mu\text{m}$, the value of d_1 was varied from 2 to $12\text{ }\mu\text{m}$, with the target thickness fixed at $6\text{ }\mu\text{m}$ across all simulations. The laser pulse simulated was linearly polarised with $I_L = 1 \times 10^{21}\text{ W/cm}^2$. The influence of this on the $2\omega_L$ conversion efficiency is shown in figure 4.24.

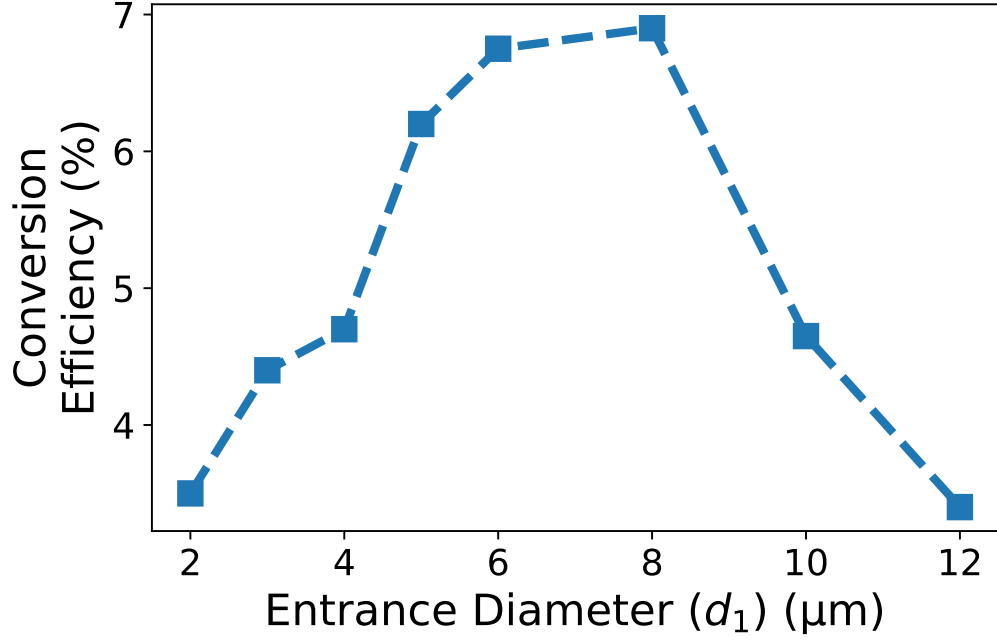


Figure 4.24: Conversion efficiency to $2\omega_L$ as a function of the entrance diameter d_1 , for a fixed exit diameter $d_2 = 2 \mu\text{m}$ and $l = 6 \mu\text{m}$ with a pulse with $I_L = 1 \times 10^{21} \text{ W/cm}^2$.

Increasing the entrance aperture diameter leads to an enhancement in conversion efficiency, reaching a peak of 6.9% when d_1 is set to $8 \mu\text{m}$. This optimal efficiency arises due to the convergence angle of the self-focusing laser pulse matching the angle of the aperture walls. The laser pulse then focuses to its highest intensity near the exit aperture, as can be observed in figure 4.25(b). The importance of the relative angle between focusing laser pulse and aperture walls on the conversion efficiency is made clearer by comparing figures 4.25(a) and (b). The former is the case of $d_1 = 3 \mu\text{m}$, for which a mismatch between the two angles leads to the focal position of the laser pulse occurring inside the target, resulting in a decrease in efficiency, similar to the case shown in figure 4.16. As the size of the entrance aperture is increased beyond $8 \mu\text{m}$, the focal position of the pulse moves further past the exit of the channel, leading to reduced interaction with the aperture walls. Further tailoring of the target geometry may enhance conversion efficiency, potentially allowing for even greater alignment between the laser's self-focusing properties and the target. Additionally, the optimum may also depend on the temporal intensity contrast of the laser pulse. However, exploring these factors lies beyond the scope of this chapter.

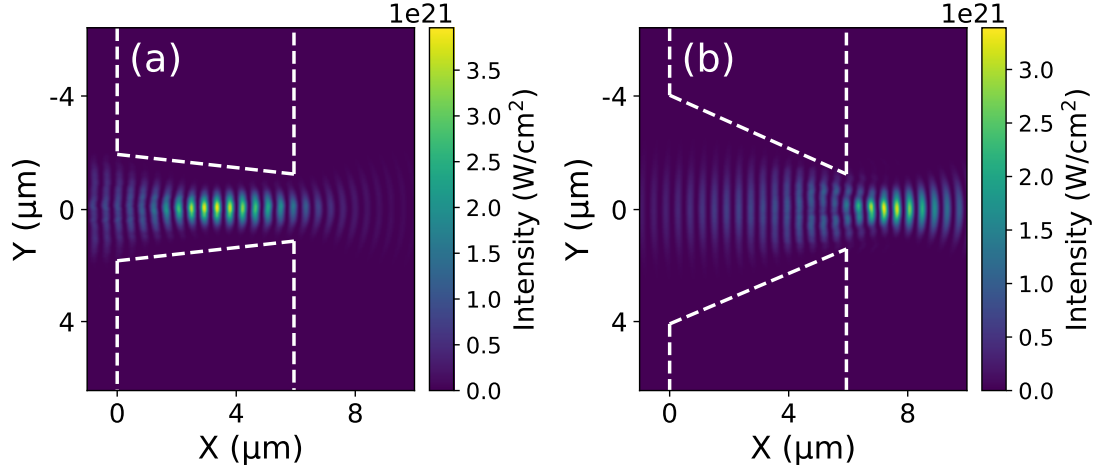


Figure 4.25: (a) Intensity profile of the drive laser light in the XY plane at $Z = 0 \mu\text{m}$, from 3D simulations, for $d_1 = 3 \mu\text{m}$ ($d_2 = 2 \mu\text{m}$ and $l = 6 \mu\text{m}$) at $t = 23 \text{ fs}$. (b) Is the same conditions as (a) for $d_1 = 8 \mu\text{m}$.

4.7 Conclusion

This chapter reports on the numerical investigation performed to explore the use of a preformed plasma aperture to generate intense second harmonic light with higher order modes. This study was built on the work previously published by M. Duff *et al* [56], in which light with similar properties was generated from relativistic plasma apertures formed during the interaction of intense laser light with an ultra-thin targets.

The results of the numerical investigation demonstrate that incoming laser light with a Gaussian spatial distribution that is incident on a preformed aperture target can be converted to relativistically intense second-harmonic radiation emission with higher-order spatial modes. The spatial structure of the mode was also found to be dependent on the polarisation of the drive laser pulse, with a linearly polarised pulse generating a TEM_{01} mode and a circularly polarised pulse generating an annular profile with OAM. The dependence of the conversion efficiency into frequency-doubled light on the interaction parameters of target thickness and aperture diameter was explored, finding an optimal diameter and target thickness at which any increase in target thickness does not lead to further gains in efficiency. This was shown to be proportional to the strength of the longitudinal electric field at the aperture edge, the quantity of laser light absorbed by the target and the percentage of the total available light that enters the aperture. This study was performed for linear polarisation in both 2D and 3D simulations, which broadly exhibited the same trends. This then allowed for a

multi-dimensional parameter space to be explored more readily using computationally efficient 2D simulations. Broadly similar trends in conversion efficiency were seen when circularly polarised light was used in 3D simulations.

For the parameter range explored, an optimal conversion efficiency of 5% and 3.5% was achieved for 2D and 3D simulations respectively. For a pulse with $\phi_L = 2.1 \mu\text{m}$ and $I_L = 1 \times 10^{21} \text{ W/cm}^2$, these figures were seen for aperture diameters of 2 and 3 μm respectively, with the conversion efficiency reaching a maximum and subsequently saturating for target thickness of $\sim 7 \mu\text{m}$ and greater. Increasing the intensity of the drive pulse leads to slight reductions in the efficiency due to the onset of relativistic transparency at the aperture edges, acting to increase the effective aperture diameter.

Differences between the results from 2D and 3D simulations can likely be attributed to dimensionality effects. In 3D, reduced interaction occurs between the aperture edges and the laser fields when the polarisation vector is not perfectly normal to the edges. Additionally, the extra degree of freedom in 3D allows accelerated electrons to escape the laser field's influence as they propagate. The use of tapered targets demonstrated the potential for enhancing optimal conversion efficiency. By aligning the taper angle of the aperture channel with the convergence angle of the self-focusing laser pulse, the interaction between the laser field and the aperture edge was increased, achieving a higher maximum conversion efficiency of 6.9%. This result highlights the potential for further increases in efficiency through additional tailoring of target geometry. Finally, an inverse relationship between the aperture diameter and the size of the generated mode structure was observed, demonstrating the ability to tune the mode size through varying target geometry.

Chapter 5

Numerical simulations to guide the experimental investigation of $2\omega_L$ light generated from aperture targets

5.1 Introduction

This chapter reports on a simulation-based study to guide the success of a planned experimental investigation into the generation of frequency doubled light from aperture targets using the Gemini laser system (discussed in section 3.3.1). The simulations reported on in the previous chapter modelled an idealised interaction between a pulse and an aperture target. Not included were several features present during laser-plasma interaction experiments that could potentially have a great impact on interaction dynamics and thus the generation of $2\omega_L$ light. The first of these is the generation of pre-plasma within the aperture channel due to the light present prior to the arrival of the main pulse (discussed in section 3.2.2). If the density of this pre-plasma in the region of the laser propagation axis approaches or exceeds critical density, it would likely impact the propagation of the pulse through the aperture. This could impact the generation of electron bunches from the aperture edges and consequently the generation of $2\omega_L$ light. The second feature is the alignment of the laser focal position to the aperture spatial location, known as the on-shot spatial jitter of the laser pointing.

If the magnitude of this movement is such that either a lower intensity portion of the pulse interacts with the aperture edge or the pulse misses the aperture target entirely, $2\omega_L$ generation will be affected.

To explore the influence of these effects, it was necessary to employ a combination of simulation types. Accurately modelling the pre-plasma generation necessitated the use of a hydrodynamic code, specifically FLASH, capable of simulating an interaction over the timescales (hundreds of ps) and pulse intensities (10^8 to 10^{14} W/cm²) over which the light preceding the main pulse exists. The results of these simulations were then used in combination with EPOCH simulations to model the effects of both the pre-plasma formation and the on-shot spatial jitter of the laser pointing on $2\omega_L$ generation.

5.2 Pre-plasma formation

As covered in section 3.2.2, high-power laser pulses are typically preceded by both a pre-pulse and an ASE pedestal. These are often of sufficient intensity ($\geq 10^{12}$ W/cm²) to cause ionisation of the target before the arrival of the main pulse, generating a pre-plasma. The influence of this pre-plasma on interaction dynamics in the case of the use of thin foil targets has been extensively studied, as covered in section 2.5. It is therefore reasonable to assume that it may also influence the generation of higher order modes (HOMs) and frequency doubled light from aperture targets. The previous chapter reported on a set of simulations performed to investigate a simple pre-plasma profile, that of a uniform filling of the aperture channel with varying densities of plasma up to the critical density. As demonstrated, the presence of this pre-plasma does influence the generation of $2\omega_L$ light from aperture targets. While this was a simplistic approach to determine the impact of plasma filling at the aperture surface, these results do justify further exploration of the role of pre-plasma on $2\omega_L$ generation.

5.2.1 Modelling the influence of pre-plasma with EPOCH

The investigation of the influence of the pre-plasma on the generation of frequency-doubled light began with the use of EPOCH simulations. While they do not model the physical processes occurring at lower pulse intensities that are responsible for the generation of pre-plasma, they do allow for simulations to be initialised with user defined plasma density distributions. The assumption was made that the pre-pulse and

ASE pedestal would cause the generation of plasma along the edges of the aperture, and that this plasma would then expand toward the aperture center, analogous to the process of ‘pinhole closure’ due to plasma formation in a spatial filter incorporating focusing through a pinhole [107]. As such, a transverse pre-plasma profile was defined within the EPOCH input deck using the following function;

$$f(y) = 1 - e^{-(y/\phi_w)^n} \quad (5.1)$$

where ϕ_w defines the dimensions of the super-Gaussian to the power n and therefore the aperture width. The extent of pre-plasma expansion into the center of the aperture channel is then varied by changing the value of n . A series of 2D EPOCH simulations were then performed with varying levels of pre-plasma expansion defined by multiplying this function by the peak target density. For these simulations, the grid was defined as 1000×720 mesh cells, corresponding to $20 \mu\text{m} \times 20 \mu\text{m}$. The targets were comprised of Al^{13+} ions, with the density defined by equation 5.1 for a fixed value of $\phi_{fwhm} = 3.6 \mu\text{m}$. This value was chosen so that the pre-plasma density profiles were similar in size to a $6 \mu\text{m}$ aperture with no pre-plasma present, a case that was also simulated. This provided a control against which the effects of varying plasma expansion into the aperture channel on $2\omega_L$ generation could be compared. A target thickness of $l = 6 \mu\text{m}$ was used. Figure 5.1(a) shows line-outs taken through the aperture target, demonstrating the pre-plasma profile generated for varying levels of n in equation 5.1, as well as the no pre-plasma case. The target ions were neutralised with an electron population with a peak density of $30n_{crit}$ and an initial temperature of 10 keV. Each species was initialised with an average of 15 particles per cell. All the boundaries of the simulation box were defined as free-space, and the ions were mobile in all cases. The pulse of laser light with wavelength $\lambda_l = 800 \text{ nm}$, matching that of the Gemini system, entered from the left (X) boundary and was focused at the plane corresponding to the front surface of the aperture ($X = 0$). Both the laser pulse temporal and spatial profiles were defined in terms of intensity as Gaussian with pulse duration of $t_L = 40 \text{ fs}$ (FWHM) and focal spot diameter of $\phi_L = 3 \mu\text{m}$ (FWHM). Time $t = 0$ was defined as when the peak of the pulse reaches $X = 0$. The peak laser pulse intensity was held constant at $I_L = 1 \times 10^{21} \text{ W/cm}^2$.

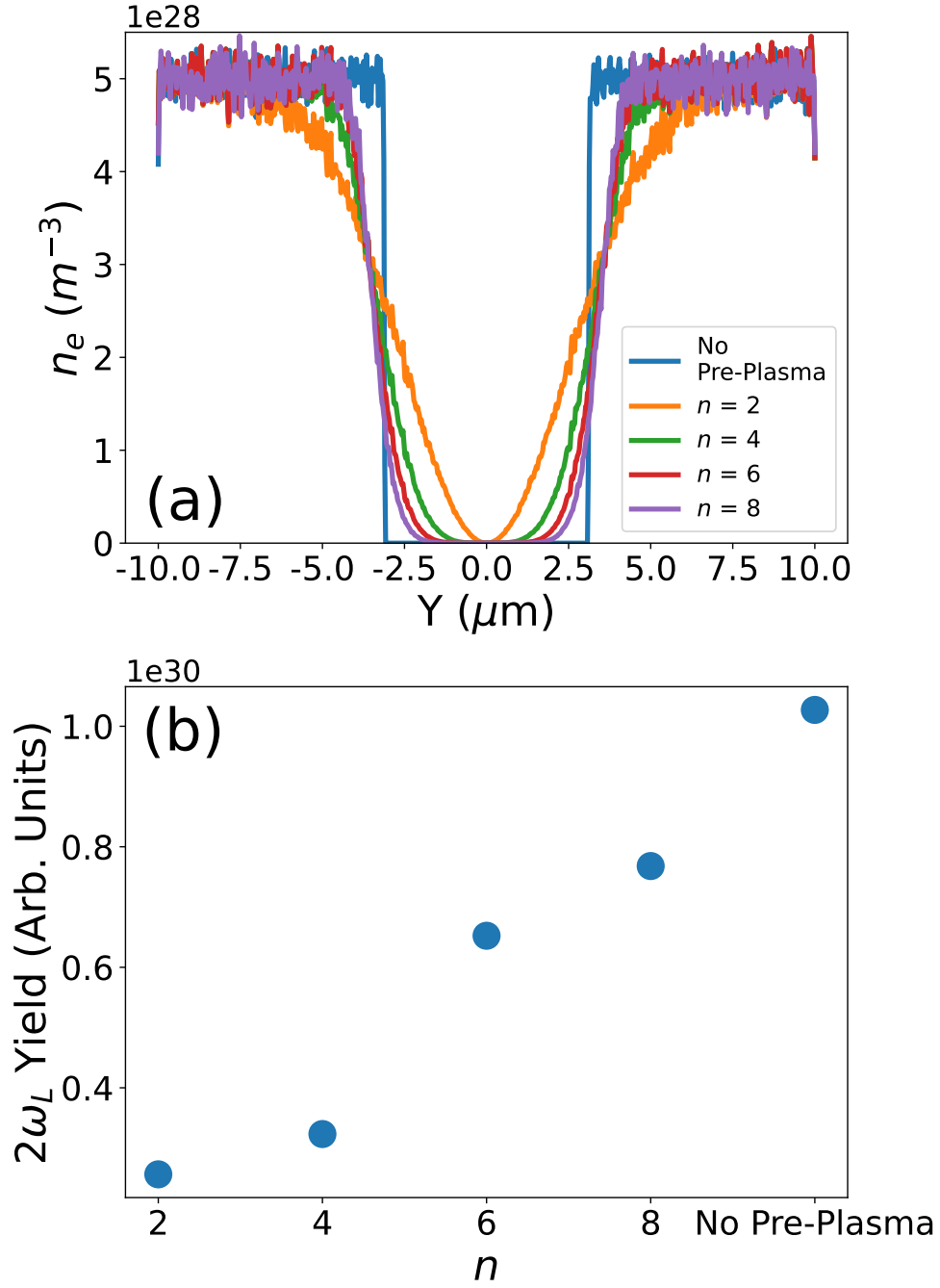


Figure 5.1: (a) Line-outs through the aperture channel from 2D EPOCH simulations performed to investigate the influence of varying levels of pre-plasma expansion on $2\omega_L$ yield, taken at start of the simulation, for varying values of n in equation 5.1. Also included is an aperture with no pre-plasma present. (b) The $2\omega_L$ yield as a function of pre-plasma expansion, taken from the 2D EPOCH simulations as a function of n .

Figure 5.1(b) shows the $2\omega_L$ yield as a function of n , corresponding to different pre-plasma profiles and the case of no pre-plasma present. As shown in the plot, the presence of pre-plasma affects $2\omega_L$ generation, with any degree of expansion reducing the yield. This reduction becomes more substantial as the degree of pre-plasma ex-

pansion increases. This change in yield is significant enough to justify the use of the hydrodynamic code FLASH, the details of which are discussed in section 3.5.2, to gain a more accurate understanding of the generation of pre-plasma.

5.2.2 Convergence testing of FLASH

To ensure that reliable results were obtained from the use of FLASH, a series of tests were performed. Initially, the modelling of the laser pulse behaviour within FLASH was tested to ensure that the temporal intensity profile and focal-spot size were correctly generated. A suitable simulation resolution and maximum adaptive mesh refinement level were then determined, both via convergence testing.

Pulse temporal intensity profile

A crucial aspect of understanding the influence of the pre-pulse and ASE pedestal on an interaction is ensuring their accurate representation within FLASH. Figure 5.2(a) shows the temporal intensity profile of the pre-pulse and ASE pedestal measured on the Gemini system, the basis of which was used for the input into the hydrodynamic simulations. Of note are the pre-pulses seen in the orange shaded region, which are thought to be generated by the diagnostic used to measure the temporal intensity profile, not representative of any actual change in the intensity at this time prior to the main pulse. Potential causes of these artefacts are discussed in [83]. As such, these were not to be included in the final FLASH simulations.

Within the FLASH code, the temporal intensity profile of a pulse is defined by a series of pulse powers and corresponding time points, with the number of time steps limited to 20. Using the temporal intensity contrast measurement made by the Gemini laser team, a series of time steps and powers could be selected. However, taking them directly from figure 5.2(a) could lead to the selection of the diagnostic-generated pre-pulses mentioned above. To avoid this, a median filter was applied, the result of which is shown in figure 5.2(b) along with the original profile. From this median filtered data, 20 time points were selected, shown in figure 5.2(c). Note that the pre-pulse that exists at ~ 10 ps prior to the main pulse is not thought to be an artefact from the diagnostic and so has been included in the selected time points.

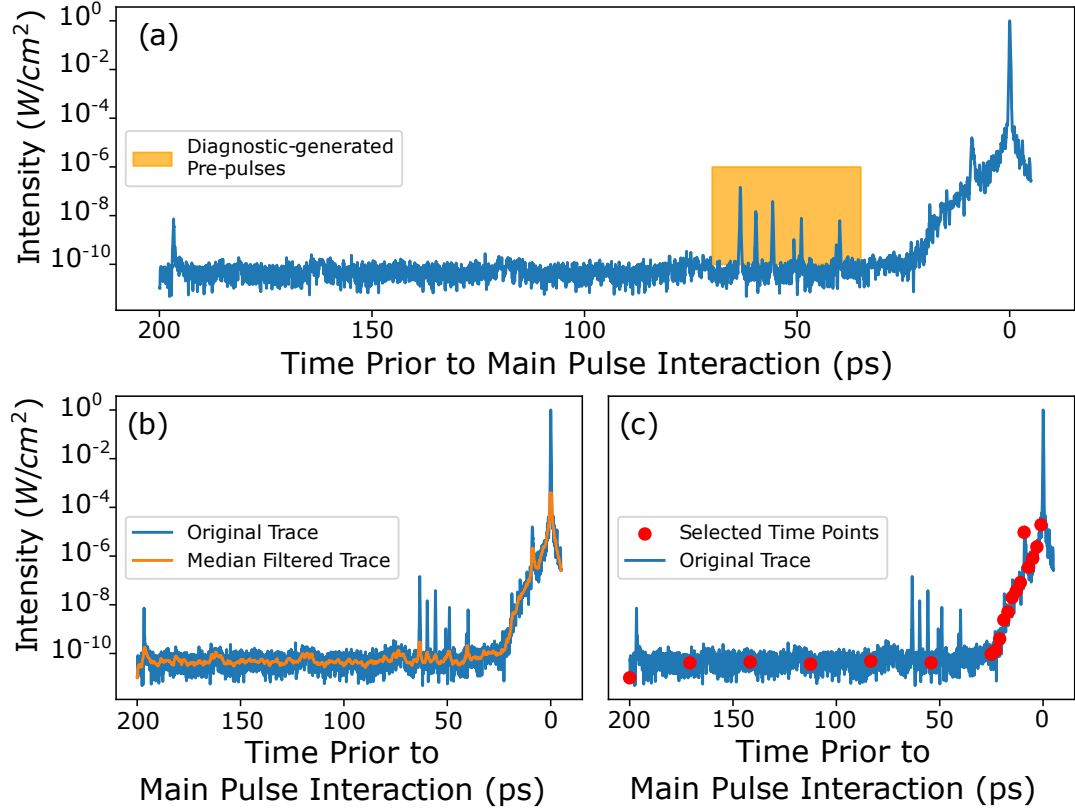


Figure 5.2: (a) A pre-pulse temporal intensity profile from the Gemini system with diagnostic-generated pre-pulses highlighted. (b) The same profile plotted along with a median filtered trace required to remove unwanted intensity spikes. (c) The same profile plotted with the time and intensity points that were selected from it for use within FLASH simulations.

Focal spot size

With an accurate temporal profile defined, the next step was ensuring that the focal spot size could be resolved. As covered in section 3.5.2, the resolution of a FLASH simulation is determined in part by the ‘block’ numbers in the X and Y dimensions. A block in FLASH is a fixed-size area of grid cells that serves as the basic unit of computation within the simulation. The focal spot size that was required to be resolved was $\phi_L = 3 \mu\text{m}$. This was a consequence of the F/2 focusing off-axis parabola (OAP) that was to be used during the experiment. The details of the experimental setup and focusing optic are covered in section 6.2 of the next chapter. Convergence testing was conducted by performing a series of simulations in which the block number in the transverse, or Y , dimension was incrementally increased. A Gaussian function was fit to the focal spot generated within each simulation to allow for its FWHM to be measured.

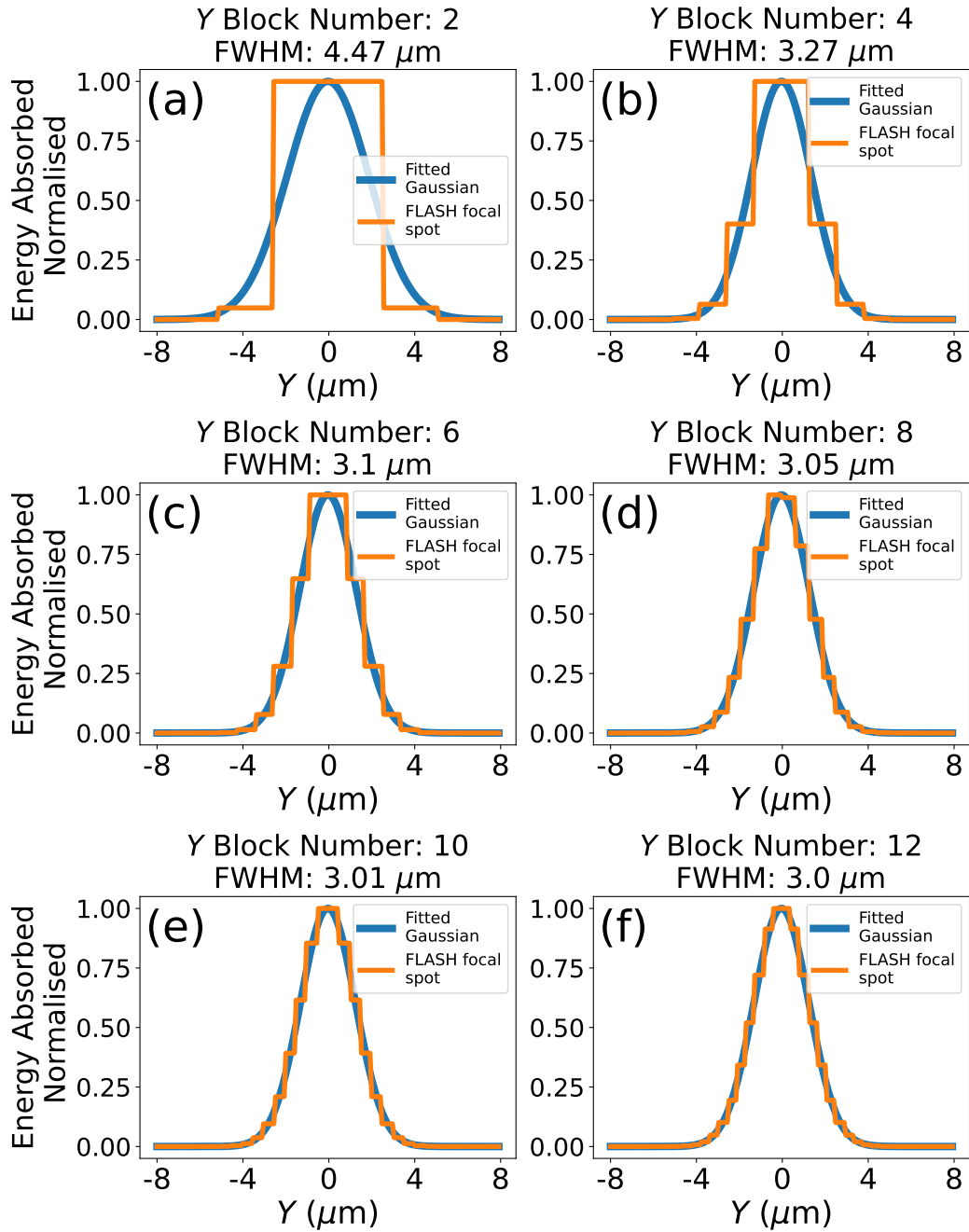


Figure 5.3: Figures showing line-outs taken from the focal spots generated in FLASH and the calculated FWHM of a Gaussian function fitted to them as the block number in the Y-dimension is increased.

Figures 5.3(a) through (f) show the focal spots, the fit Gaussian functions and the measured FWHM for each block number used in this set of simulations. Note that FLASH does not simulate the EM fields of the laser pulse, instead simulating light rays that deposit their energy in the material they are propagating through, hence the units on the y -axis in figure 5.3 are that of absorbed energy. As can be seen, a lower block number results in a sub-optimal focal spot shape. As the block number is increased,

the lineout becomes increasingly Gaussian, with the final block number of 12, shown in figure 5.3(f), allowing for a focal spot with a FWHM of $3 \mu\text{m}$ to be resolved.

Resolution

With the various pulse parameters now tested, the next stage was to determine a suitable resolution for the simulation box via convergence testing. To do this, a series of simulations were performed in which a pulse, as defined by the trace shown in figure 5.2, interacts with a solid density target of thickness $6 \mu\text{m}$. The block number in each dimension was incrementally increased and the resulting target expansion was examined by taking line-outs through the expanded target along each axis.

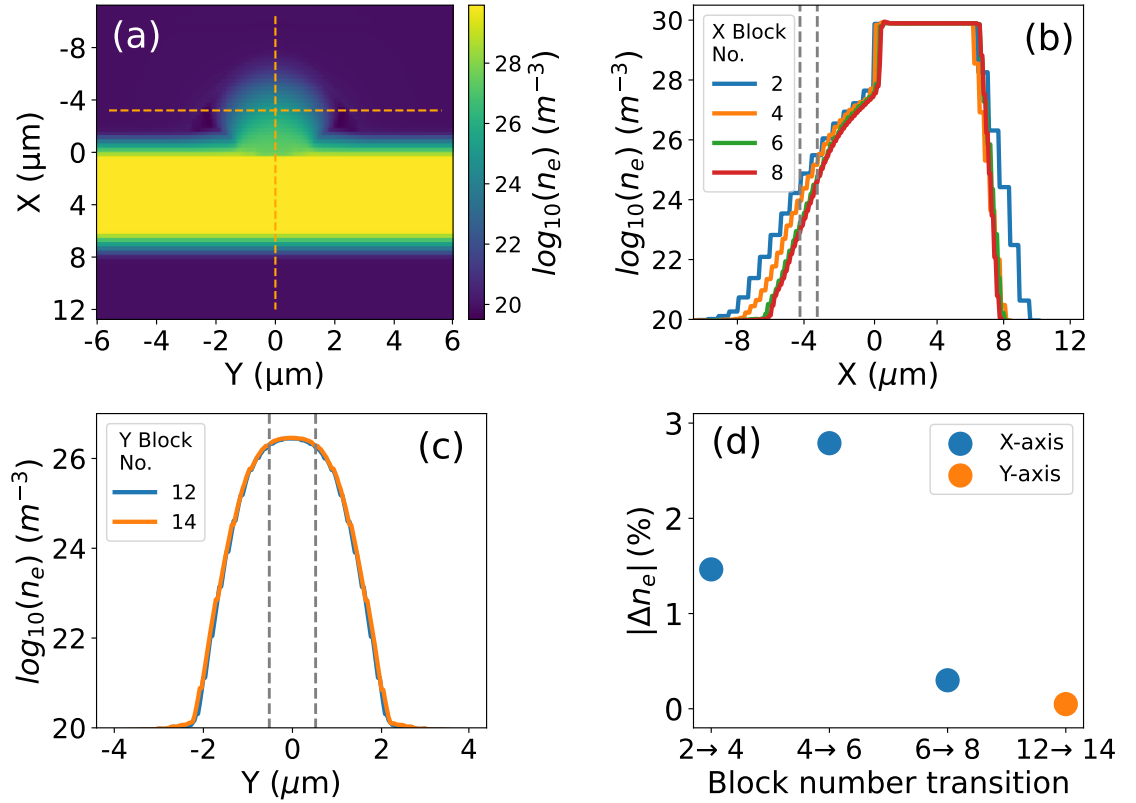


Figure 5.4: (a) Target expansion due to the interaction of a pulse with the temporal intensity profile shown in figure 5.2 for a block number of 6 and 12 in the X - and Y -dimension respectively. Also shown are position of the line-outs displayed in figures (b) and (c). (b) Electron density along the X -dimension at $Y = 0 \mu\text{m}$ for varying block number. (c) Electron density along the Y -dimension at $X = -3 \mu\text{m}$ for varying block number. (d) Absolute percentage change in average electron density averaged over a $1 \mu\text{m}$ distance indicated by the grey dashed lines from (b) and (c), plotted as a function of block number transition.

Figure 5.4(a) shows an example of target expansion for block numbers of 6 and 12 in the X - and Y -dimension respectively, with the position of the line-outs taken to measure

the expansion shown by the orange dashed lines at the positions of $X = -3 \mu\text{m}$ and $Y = 0$. The change in the degree of target expansion with increasing block number in the X - and Y -dimensions are shown in figures 5.4(b) and (c), respectively. As can be seen in figure 5.4(b), as the block number in the X -dimension is increased from 6 to 8, minimal change is seen in the target expansion. From the tests performed to ensure that the focal spot could be resolved, a minimum block number in the Y -dimension had been set at 12, but one further simulation was performed in which this number was increased to 14. In this case, minimal difference was seen, as can be observed in figure 5.4(c).

To quantify the changes in target expansion with increasing block number, the electron density was averaged over $1 \mu\text{m}$ in both the X and Y dimensions. The position of this $1 \mu\text{m}$ extent is indicated by the grey dashed lines in figures 5.4(b) and (c). Figure 5.4(d) shows the absolute percentage change in average electron density as a function of block number transition in both the X - and Y -dimensions. As can be seen in this figure, an increase in block number from 6 to 8 results in only a minimal change in electron density of approximately 0.3%. The same is true in the Y -dimension as an increase in the block number from 12 to 14 results in a change in electron density of approximately $\sim 0.05\%$. As such, block numbers of 6 and 12 were chosen for the X - and Y - dimensions respectively.

Maximum refinement level of the adaptive mesh resolution

The final convergence testing that was performed was that of increasing the maximum refinement level (MRL) for the adaptive mesh resolution (AMR) feature within FLASH, with the details of this feature covered in section 3.5.2. Using the block numbers converged upon in the previous section, the maximum refinement level was increased over four simulations from 2 to 5, with the results shown in figures 5.5(a) and (b). Note that the line-outs are sampled at the same X and Y positions as in figure 5.4(a).

As can be seen, the degree of target expansion does not appear to vary greatly with increasing MRL. To ensure that this is the case the electron density was averaged over a $1 \mu\text{m}$ distance for each dimension. The $1 \mu\text{m}$ extent is indicated by the grey dashed lines in figures 5.5(a) and (b). The absolute percentage change in averaged electron density is plotted as a function of MRL transition in figure 5.5(c). This figure demonstrates that minimal variation is observed as the MRL is increased beyond 4,

with a 0.5% and 0.07% change observed in the X and Y dimensions respectively. As such, an MRL value of 4 was chosen for the simulations performed to investigate the pre-plasma expansion.

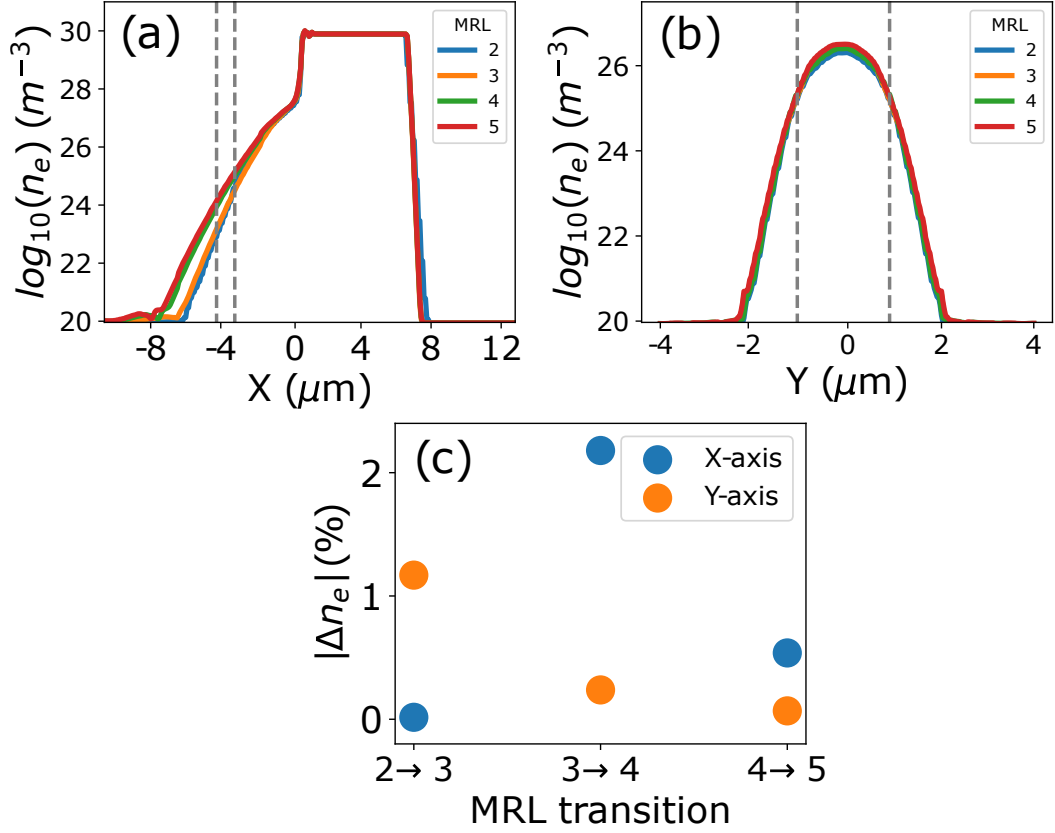


Figure 5.5: (a) and (b), the change in target expansion with varying MRL in the x- and y-dimension respectively. (c) Absolute percentage change in electron densities averaged over a $1 \mu m$ distance indicated by the grey dashed lines from (a) and (b), plotted as a function of MRL transition.

5.2.3 Modelling the formation of pre-plasma within an aperture

With the convergence testing completed, the generation of pre-plasma within the aperture channel was modelled using FLASH simulations. This began with an investigation into the influence of the intensity of the light preceding the main pulse. This involved varying the intensity of the ASE/pre-pulse portion of the intensity profile in figure 5.2 by a factor, f_I , over the range 0.01 to 10. The choice to investigate values of $f_I < 1$ was based on the potential use of plasma mirrors in the beam-line during the experiment, thereby reducing the intensity of the light preceding the main pulse by several orders of magnitude. Values of $f_I > 1$ were investigated to test the behaviour of the $2\omega_L$ generation in the case of artificially increased contrast, for example through the

introduction of a lower intensity pulse prior to the main pulse. These FLASH simulations were performed using the input parameters decided upon in the previous section. Target dimensions of $d = 6 \mu\text{m}$ and $l = 6 \mu\text{m}$ were used. While this is a larger aperture diameter than the optimal reported on in the previous chapter, this was chosen based on the knowledge of the targetry that would be available during the planned experimental campaign, further discussion of which is present in the next chapter. Figure 5.6(a) through (c) show the pre-plasma profiles generated from FLASH for the cases of $f_I = 0.01$, 1 and 10, respectively. Figure 5.6(d) shows line-outs taken through the aperture channels for each value of f_I used.

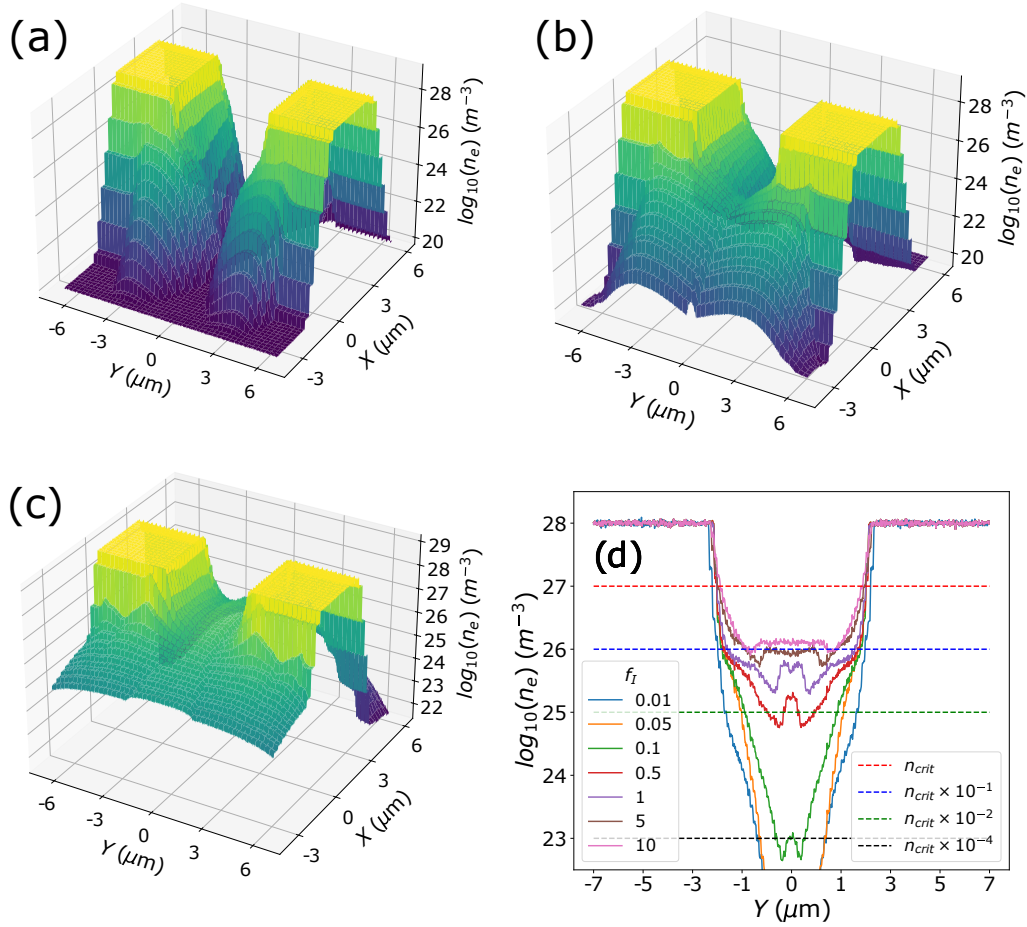


Figure 5.6: Results from FLASH simulations modelling the generation of pre-plasma from the temporal profile shown in figure 5.2 for a target with $d = 6 \mu\text{m}$ and $l = 6 \mu\text{m}$ for (a) $f_I = 0.01$, (b) $f_I = 1$ and (c) $f_I = 10$. (d) Electron density through the center of the pre-plasma profiles for all values of f_I used. The dashed red line indicates n_{crit} .

As can be seen from these plots, the extent of plasma expansion from the aperture edges increases with increasing f_I . While the density of the generated plasma does not

exceed n_{crit} , it does approach $0.1 \times n_{crit}$ for $f_I \geq 1$. At these densities, the propagation of the main pulse through the aperture channel would be altered due to the change in the plasma refractive index, impacting the generation of electron bunches from the aperture edges and therefore $2\omega_L$ light and HOM generation. While narrowing of the aperture channel occurs at all densities, the degree of narrowing is minimal for all values of f_I at $n_e \approx n_{crit}$, with plasma expanding into the channel by only around $0.5 \mu\text{m}$ from each edge. At $n_e = n_{crit} \times 10^{-1}$, the aperture channel has closed for f_I values of 5 and 10, while for the other values of f_I , the narrowing remains $\approx 0.5 \mu\text{m}$. At $n_e = n_{crit} \times 10^{-2}$, the aperture has closed for f_I values of 0.5 and 1. For $f_I = 0.01$, the channel has narrowed by $0.6 \mu\text{m}$, and for $f_I = 0.05$ and 0.1 , it narrows by $\approx 1.35 \mu\text{m}$. Finally, at $n_e = n_{crit} \times 10^{-4}$, the aperture has closed for $f_I = 0.1$, while for the remaining values of f_I , the channel narrows by $\approx 2 \mu\text{m}$. Given the importance of the aperture diameter to focal spot size ratio on $2\omega_L$ generation, as determined from the simulations performed in the previous chapter, this narrowing could result in a change in the optimal focal spot size. These results justified further investigation of the impact of pre-plasma on the interaction through the use of EPOCH simulations.

5.2.4 Modelling the influence of pre-plasma on $2\omega_L$ generation

The pre-plasma density distributions generated from FLASH for each value of f_I were used as the initial particle density distributions for a set of EPOCH simulations. For brevity, this approach will hereafter be referred to as ‘combined FLASH and EPOCH simulations’. A particle-per-cell number of 15 was used for a simulation box of 3200×3200 mesh cells, corresponding to $45 \mu\text{m} \times 45 \mu\text{m}$. A 40 fs pulse with a focal spot size of $\phi_L = 3 \mu\text{m}$ (FWHM) and pulse intensity of $I_L = 1 \times 10^{21} \text{ W/cm}^2$ was used, with a total simulation time of 200 fs. Figure 5.7 shows the change in the $2\omega_L$ yield for all values of f_I . Also included is the case of no pre-plasma for comparison.

As can be seen, the presence of a pre-plasma within the aperture channel has a significant effect on the $2\omega_L$ yield. When examining the simulation in which no pre-plasma is present, it can be seen that the $2\omega_L$ yield is increased in all cases excluding $f_I = 10$, for which is approximately the same. An optimal pre-plasma is observed for $f_I = 0.1$, for which the $2\omega_L$ yield is increased by $\approx 3.3\times$ in comparison to the no pre-plasma case.

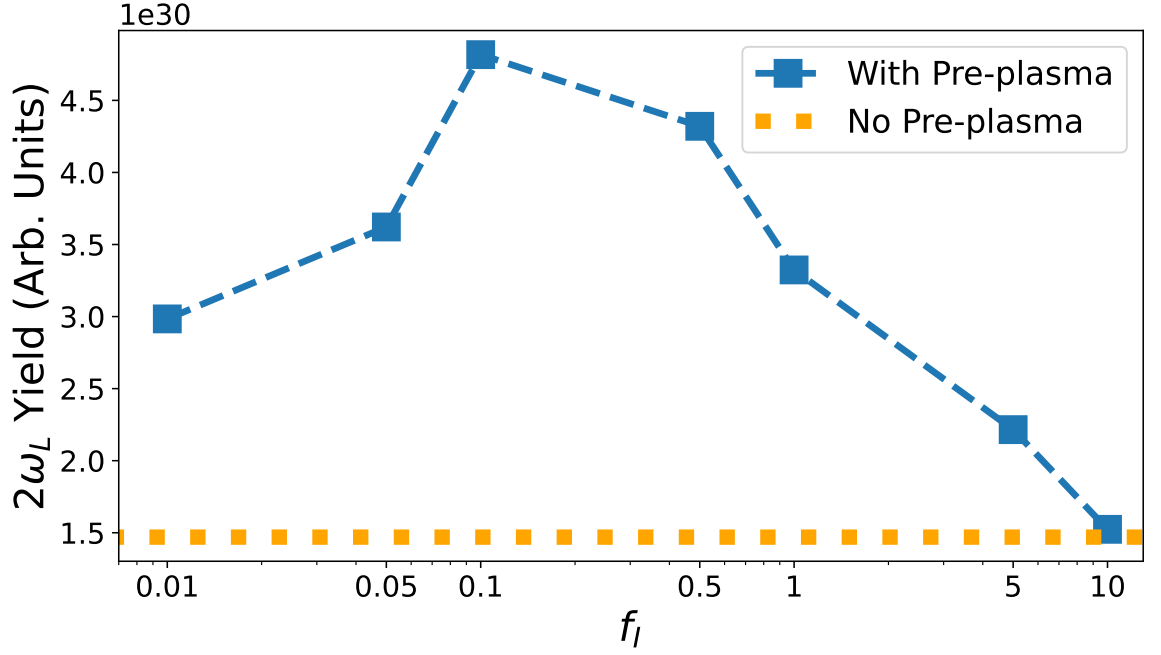


Figure 5.7: $2\omega_L$ yield as a function of f_I , taken from combined FLASH and EPOCH simulations for a target with $d = 6 \mu\text{m}$ and $l = 6 \mu\text{m}$.

The increase in $2\omega_L$ yield observed between $f_I = 0.01$ and 0.1 is likely to be due to the narrowing of the aperture channel. This narrowing leads to a higher intensity portion of the pulse interacting with the aperture edges, thus generating higher density electron bunches. As f_I increases further, the filling of the aperture with plasma may begin to affect the propagation of the pulse through the channel and thus the electron bunch generation. To investigate this, the $2\omega_L$ yield as a function of simulation time is plotted for the range of f_I values investigated, shown in figure 5.8. Also shown is the squared E_Y field of the laser integrated over a $1 \mu\text{m} \times 1 \mu\text{m}$ region at $X = 0$ within the simulation box, from a simulation in which no target is present, modelling the case of the propagation of the pulse not being impeded at all. In the case of $f_I = 0.01$ and 0.1 , the peak in $2\omega_L$ generation coincides with the peak in laser field strength at an elapsed simulation time of 160 fs . This suggests that for these pre-pulse intensity levels, the propagation of the pulse within the aperture channel is not impeded.

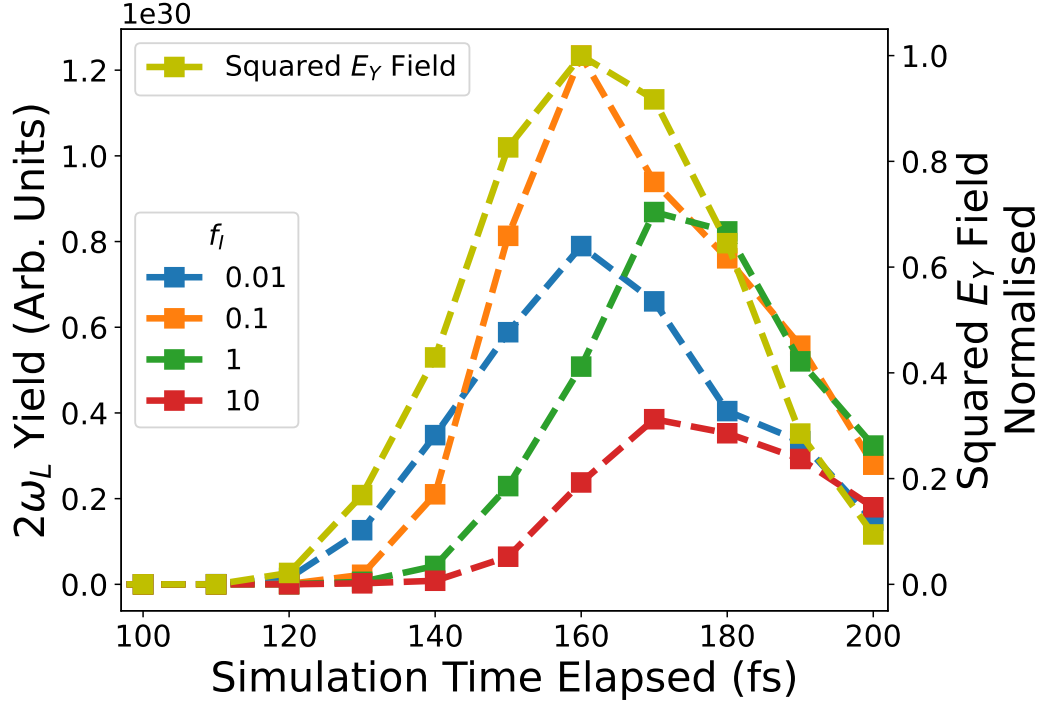


Figure 5.8: Comparison of the $2\omega_L$ yield over the course of an EPOCH simulation with the change in squared E_Y field of the laser.

For the cases of nominal intensity and $f_I = 10$, peak $2\omega_L$ generation is delayed relative to the peak in laser field strength, with both reaching a maximum value at an elapsed simulation time of 180 fs. This suggests that the main peak of the pulse is being hindered by the pre-plasma, suppressing the generation of frequency doubled light at this point during the interaction, leading to the comparatively reduced yield seen in figure 5.7.

Figures 5.9(a) through (d) show line-outs taken through the time integrated $2\omega_L$ transverse spatial distributions from the combined FLASH and EPOCH simulations for f_I values of 0.01, 0.1, 1 and 10 respectively. These demonstrate that the spatial distribution is largely maintained under all pre-pulse intensities, with the addition of lower intensity lobes present either side of the main mode. However, a more pronounced difference is observed in the spatial distribution generated for $f_I = 10$. In this case, the two central lobes of light exhibit a greater degree of asymmetry, while lobes on either side are closer in intensity to those in the center, indicating that increased amounts of pre-plasma can reduce the stability of the generation mechanism.

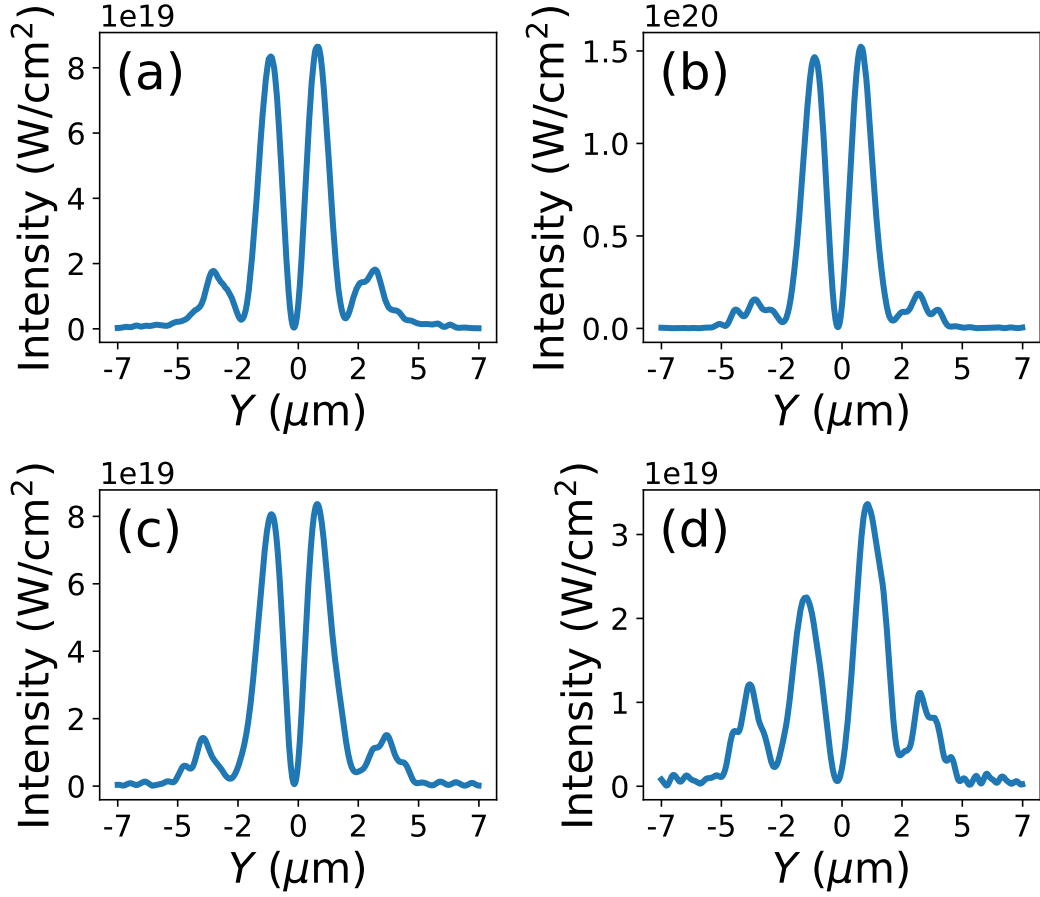


Figure 5.9: Line-outs taken through the time integrated $2\omega_L$ transverse spatial distributions from combined FLASH and EPOCH simulations for a target with $d = 6 \mu\text{m}$ and $l = 6 \mu\text{m}$ for $f_I =$ (a) 0.01, (b) 0.1, (c) 1 and (d) 10, measured $10 \mu\text{m}$ from the rear of the target.

The simulation results reported in this section are largely encouraging for the experimental campaign. They indicate that for the pre-pulse and ASE pedestal at the intensity expected to occur on the Gemini system, with the planned use of a single plasma mirror, generation of $2\omega_L$ light in the desired mode should still occur. While this generation could be somewhat enhanced through manipulation of the intensity of the light preceding the pulse, not doing so should not prevent the generation of $2\omega_L$ light in the experimental campaign.

5.2.5 Optimal spot size to aperture diameter ratio

A key finding from the simulations reported in the previous chapter is that optimal conversion of laser pulse energy into $2\omega_L$ light occurs when the focal spot size of the pulse is approximately equal to the diameter of the aperture target. Since modelling of pre-plasma formation has shown that a narrowing of the aperture channel occurs, the

optimal ratio may differ from that identified previously. To investigate this, a set of combined FLASH and EPOCH simulations were performed at nominal pre-pulse and ASE pedestal intensity in which, with a fixed spot size of $\phi_L = 3 \mu\text{m}$, the aperture size was varied from $3 \mu\text{m}$ to $6 \mu\text{m}$. Figure 5.10 shows the results of this set of simulations.

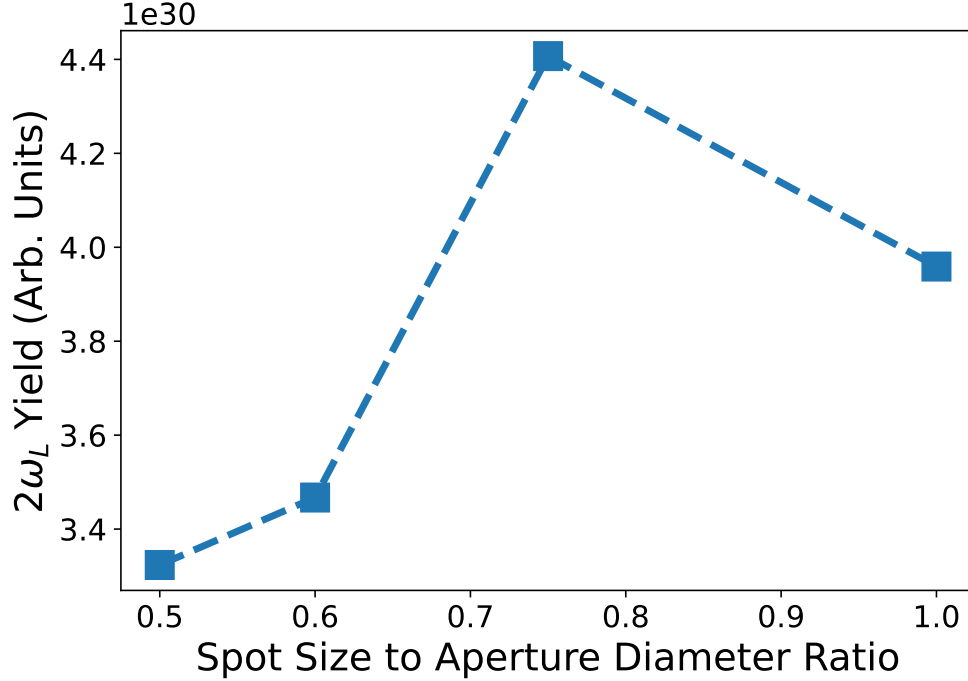


Figure 5.10: $2\omega_L$ yield as a function of spot size to aperture size ratio for fixed focal spot size $\phi_L = 3 \mu\text{m}$ and a target with $d = 6 \mu\text{m}$, taken from combined FLASH and EPOCH simulations.

This figure demonstrates that the inclusion of the pre-pulse has the effect of reducing the optimal ratio below 1, with its value most likely in the range of 0.7 to 0.9. Given that the smallest aperture diameter target available for use during the experiment is $6 \mu\text{m}$ and the focal spot size at best focus is $3 \mu\text{m}$, it is clear that the optimal focal spot size to aperture diameter ratio cannot be achieved without adjustments to the experimental set up. While deviating from this optimal ratio does not completely halt $2\omega_L$ light generation, the results of these simulations make clear the fact that an increase in $2\omega_L$ yield is achievable. One possible method of attaining a larger focal spot size is to move the target away from best focus, closer to the focusing optic. This technique, referred to as defocusing, increases the focal spot size at the point of interaction. However, this is accompanied by a decrease in pulse intensity, which would in turn reduce the potential $2\omega_L$ yield.

5.3 On-shot spatial jitter of the laser pointing

Prior to a full power shot, the position of the laser focal spot, known as the laser pointing, must be aligned with the centre of the aperture target. This alignment is achieved using a low-power continuous-wave (CW) beam that is propagated through the same chain of optics as the full-power pulse. This ensures that the focal spots of the low-power CW beam and full power beam are overlapped. However, due to factors such as mechanical vibrations of the optics that make up the laser chain [108][109], the position of the CW beam focal spot and that of the full-power pulse can vary between the time that the alignment process has been completed and the full-power shot is taken. This phenomenon, known as the on-shot spatial jitter of the laser pointing, or spatial jitter for convenience, is inherent in the Gemini laser system and all other high-power laser systems. Depending on the magnitude of the spatial jitter, relative to the focal spot size and aperture diameter, there are several ways in which the generation of $2\omega_L$ light could be negatively impacted. A magnitude that results in the centre of the focal spot being within, but not centred on, the aperture could lead to a lower intensity portion of the pulse interacting with the aperture edges, or only one edge of the target. This could then reduce electron bunch generation and potentially alter the spatial properties of the generated light. Spatial jitter of a magnitude such that the focal-spot of the laser pulse interacts with only the bulk of the target could lead to a complete halting of $2\omega_L$ light generation.

To gain an understanding of the influence of this phenomenon on the interaction, statistical spatial jitter measurements were made using the Gemini system prior to the experiment, shown as a histogram in figure 5.11. These measurements consist of 300 shots taken in succession without any adjustments made to the position of the focal spot between each shot. The magnitude of the radial distance from the focal spot captured for each shot to the initial focal spot was then measured. The laser was used in a low-power mode when these measurements were made, allowing for the direct imaging of the focal spot that would not be possible were the system operating in a high-power mode. This data consists of 300 shots taken in succession without any adjustments made to the position of the focal spot between each shot. These movements are due to vibrations of the optics within the laser chain and can therefore be said to be a good representation of the actual spatial jitter that would occur during the experiment.

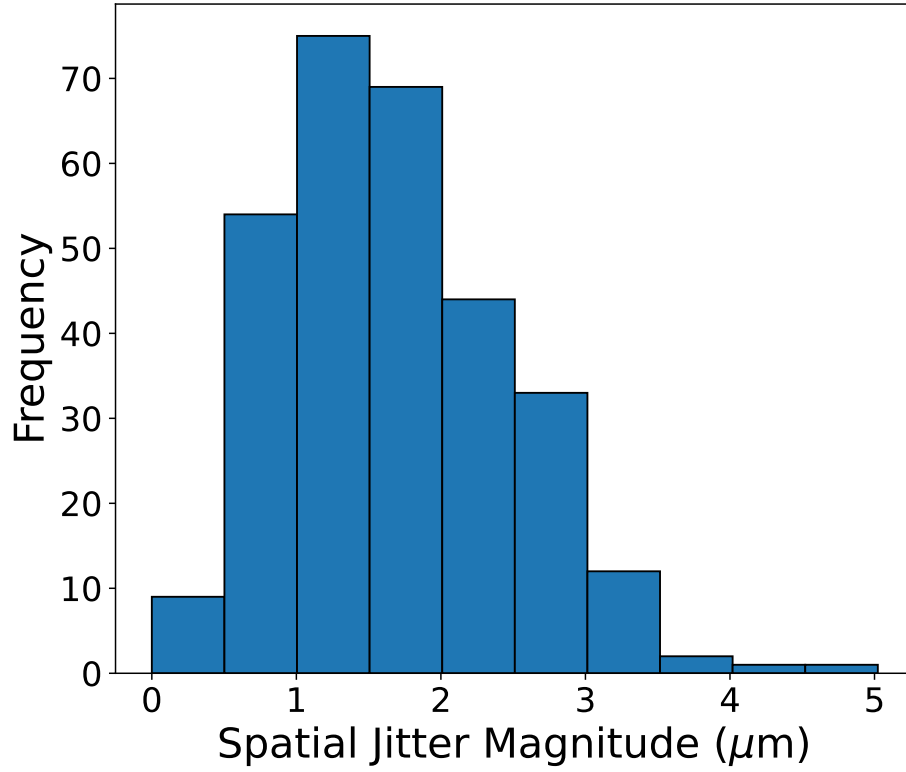


Figure 5.11: Histogram displaying the frequency of spatial jitter magnitude over 300 shots, taken using a low-power mode of the Gemini Laser System.

These measurements demonstrate that the most likely spatial jitter magnitude is $1.5 \mu\text{m}$, and the least likely is $4.5 \mu\text{m}$. The probability of a spatial jitter magnitude of $0 \mu\text{m}$ occurring is only $\sim 3\%$, while there is a $\sim 63\%$ probability that the jitter lies between 1 and $2.5 \mu\text{m}$. As this range is comparable in size to both the $6 \mu\text{m}$ aperture diameter of the targets that will be used throughout the experiment and the $3 \mu\text{m}$ focal spot size at best-focus, this range of spatial jitter magnitude would be likely to affect $2\omega_L$ generation from the target. This high likelihood of non-negligible spatial jitter therefore justifies further investigation into its effect on the interaction.

5.3.1 Modelling the influence of spatial jitter on $2\omega_L$ generation

To model the influence of the spatial jitter on the generation of both $2\omega_L$ light and the spatial mode shape, a set of combined FLASH and EPOCH simulations were performed in which the position of the laser in Y -direction was incrementally varied from the center of the aperture to a maximum displacement of $5 \mu\text{m}$. This range was chosen to match the greatest spatial jitter magnitude measured prior to the experiment, as shown in figure 5.11. This set of simulations was performed with $f_I = 1$ and for two different

aperture diameters, $d = 6 \mu\text{m}$ and $10 \mu\text{m}$. These dimensions were chosen to match the two aperture diameters that would be available during the experiment. In all other respects, the simulation parameters remained the same as those reported earlier in this chapter that modelled the influence of the pre-plasma on $2\omega_L$ generation. The results of these simulations are shown in figure 5.12 in which the $2\omega_L$ yield as a function of spatial jitter for the two aperture diameters is plotted.

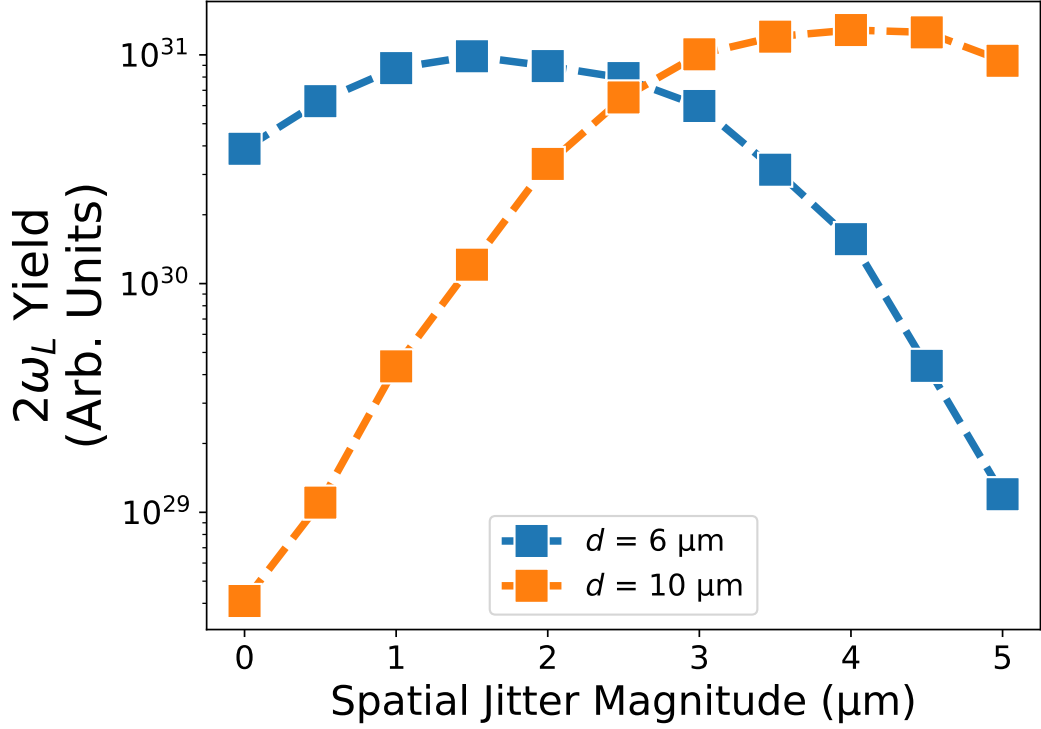


Figure 5.12: The $2\omega_L$ yield as a function of spatial jitter magnitude for targets with $d = 6$ and $10 \mu\text{m}$ taken from combined FLASH and EPOCH simulations.

The results of the simulations show a significant impact on the $2\omega_L$ yield as spatial jitter is varied. For the case of the $d = 6 \mu\text{m}$ target, as the magnitude increases from 0 to $1.5 \mu\text{m}$, an increase in the yield is observed. This can be attributed to a higher intensity portion of the pulse interacting with the aperture edge. This is evident when comparing figures 5.13(a) and (b), which show the laser pulse position relative to the aperture edges for spatial jitter magnitude of $0 \mu\text{m}$ and $1.5 \mu\text{m}$ respectively. Further increases in the spatial jitter magnitude lead to the highest intensity portion of the pulse moving past the aperture edge and interacting with the bulk of the target, as shown in 5.13(c), displaying the pulse position for a spatial jitter magnitude of $3.5 \mu\text{m}$. The behaviour observed for the $d = 10 \mu\text{m}$ target is substantially different. The yield

shows a steady increase between spatial jitter magnitudes of 0 and $3.5 \mu\text{m}$. This can be attributed to the fact that focal spot size is much smaller than the aperture diameter in this case. In the absence of any spatial jitter, only the lower intensity edges of the pulse are interacting with the aperture edges, as seen in figure 5.13(d). As the spatial jitter magnitude increases, a higher intensity portion of the pulse begins to interact with the aperture edge, as shown in figures 5.13(e) and (f). A decrease in the yield is observed as the spatial jitter magnitude increases beyond $4.5 \mu\text{m}$. As with the case of a $d = 6 \mu\text{m}$ target, this can be attributed to the highest intensity portion of the pulse beginning to move past the aperture edge, interacting with the bulk of the target.

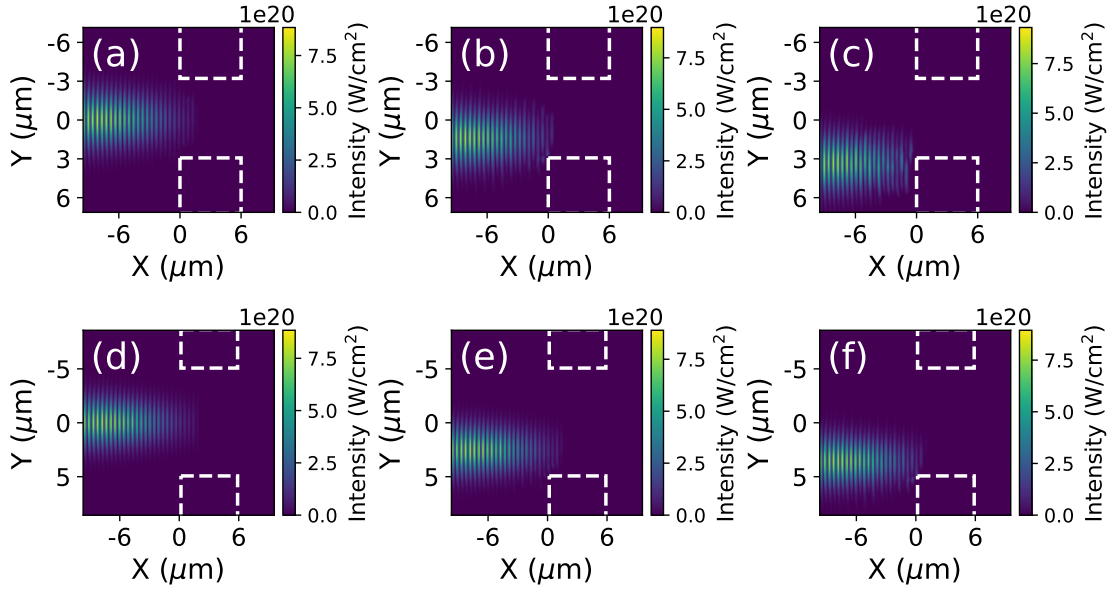


Figure 5.13: Position of the laser pulse relative to the aperture targets. (a) through (c) correspond to a target with $d = 6 \mu\text{m}$ and spatial jitter magnitudes of $0 \mu\text{m}$, $1.5 \mu\text{m}$, and $3.5 \mu\text{m}$, respectively. (d) through (f) show the equivalent results for a target with $d = 10 \mu\text{m}$, again for spatial jitter magnitudes of $0 \mu\text{m}$, $2.5 \mu\text{m}$, and $3.5 \mu\text{m}$. All figures are taken from combined FLASH and EPOCH simulations at $t = 0 \text{ fs}$, performed to investigate the influence of spatial jitter magnitude on $2\omega_L$ yield. The dashed white lines indicate the initial target profile and the nominal focused intensity of $I_L = 1 \times 10^{21} \text{ W/cm}^2$ was used.

To investigate the impact of spatial jitter magnitude on the spatial properties of the generated $2\omega_L$ light, line-outs taken through the time integrated $2\omega_L$ transverse spatial distribution for a selection of spatial jitter magnitudes are shown in figure 5.14 for both $d = 6$ and $10 \mu\text{m}$ targets. Figures 5.14(a) through (c) show the case of a $d = 6 \mu\text{m}$ target for spatial jitter magnitudes of $0.5 \mu\text{m}$, $1.5 \mu\text{m}$ and $2.5 \mu\text{m}$, respectively. For a spatial jitter magnitude of $0.5 \mu\text{m}$, the mode is largely unchanged, with two distinct lobes of light present. As the spatial jitter magnitude increases, the mode becomes

increasingly asymmetric, as electron bunches are primarily accelerated from one of the aperture edges.

Figures 5.14(d) through (f) show the case of a $d = 10 \mu\text{m}$ target for spatial jitter magnitudes of $0.5 \mu\text{m}$, $1.5 \mu\text{m}$ and $2.5 \mu\text{m}$, respectively. For the magnitude of $0.5 \mu\text{m}$, the generated light is already very asymmetric, displaying only a minor second lobe of light. For a magnitude of $1.5 \mu\text{m}$ and greater, the spatial profile appears to have reached a maximum asymmetry, as the pulse interacts with only one aperture edge. This can be observed in figures 5.13(e) and (f). These results demonstrate that spatial jitter magnitude can have a large impact on both the quality of the generated mode and $2\omega_L$ yield, depending on the focal spot size to aperture diameter ratio.

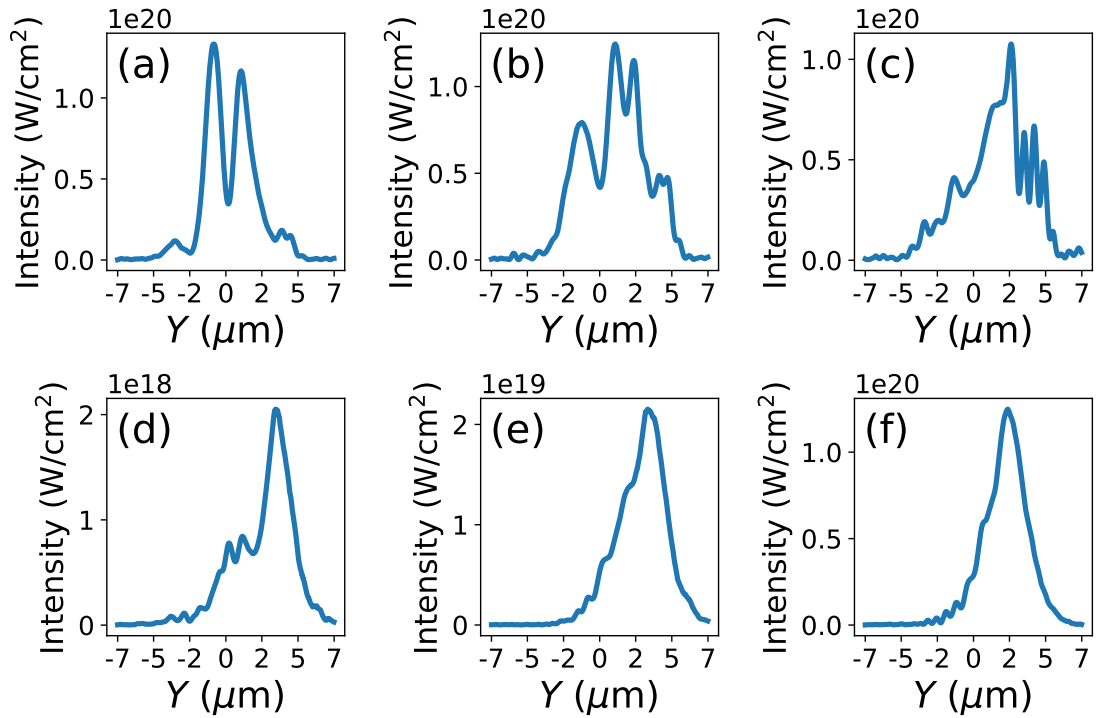


Figure 5.14: Line-outs taken through the time integrated $2\omega_L$ transverse spatial distribution measured $10 \mu\text{m}$ from the rear of the target. (a) through (c) correspond to a target with $d = 6 \mu\text{m}$ and spatial jitter magnitudes of $0.5 \mu\text{m}$, $1.5 \mu\text{m}$, and $2.5 \mu\text{m}$, respectively. (d) through (f) show the equivalent results for a target with $d = 10 \mu\text{m}$, again for spatial jitter magnitudes of $0.5 \mu\text{m}$, $1.5 \mu\text{m}$, and $2.5 \mu\text{m}$. All figures are taken from combined FLASH and EPOCH simulations performed to investigate the influence of spatial jitter magnitude on $2\omega_L$ yield.

One solution to this problem is to take a statistical approach, whereby a large number of shots are taken such that a subset will occur with a desired spatial jitter magnitude. For a $d = 6 \mu\text{m}$ target, a spatial jitter magnitude of $\leq 1 \mu\text{m}$ would likely result in the generation of the desired spatial mode while maintaining a high yield of $2\omega_L$ light.

Based on the data presented in 5.11, a spatial jitter of $1\ \mu\text{m}$ has a $\approx 20\%$ chance of occurring. This approach ensures that, while not every shot will occur under the desired conditions, a large enough number of shots would lead to a subset of high-quality data that can be analysed. However, the dependence of this approach on taking a larger number of shots leaves it susceptible to issues that routinely occur during experimental campaigns such as laser-system down time that reduces the total number of shots that can be taken over the course of an experiment.

An alternative solution is to defocus the laser to create a larger focal spot size at the point of interaction. This technique, used with the $d = 10\ \mu\text{m}$ targets, would result in the spatial jitter magnitude being smaller relative to the spot size. This could significantly mitigate the negative effects of the spatial jitter on the generated spatial mode while also achieving an optimal spot size to aperture diameter ratio as covered in the previous section. However, as previously mentioned, defocusing the laser is accompanied by a reduction in intensity that would lead to a reduction in $2\omega_L$ generation, but this compromise could be justified if the negative impact of the spatial jitter is reduced.

5.4 Conclusion

In this chapter the potential challenges involved with the experimental investigation of $2\omega_L$ light generation from aperture targets were modelled using numerical simulations, with the aim of developing methodologies to address these. Factors that could potentially influence the interaction dynamics are the generation of pre-plasma within the aperture in the region of the laser propagation axis due to the ASE pedestal and pre-pulses that precede the main pulse and the on-shot spatial jitter of the laser pointing.

This chapter began by reporting on an investigation into the effect of pre-plasma on the generation of $2\omega_L$ light from aperture targets. Initially performed using EPOCH simulations, the results demonstrated a reduction in the $2\omega_L$ yield that justified the use of a more accurate but complex simulation type, namely FLASH, a hydrodynamic code. The convergence testing of FLASH was discussed, followed by the results from its use in modelling the generation of pre-plasma for a range of ASE pedestal and pre-pulse intensities. These results demonstrate a narrowing of the aperture channel for lower intensities and a filling of the channel for higher intensities, with densities

approaching $0.1 \times n_{crit}$. This then justified the use of combined FLASH and EPOCH simulations to investigate the influence of the generated pre-plasma on both the $2\omega_L$ yield and the generated spatial mode. These simulations demonstrate that for the planned experimental setup, the desired mode is still generated and with an acceptable $2\omega_L$ yield. Further combined FLASH and EPOCH simulations were then performed to investigate a potential change in the optimal focal spot size to aperture diameter ratio, a key finding from the previous chapter. These simulations demonstrate that when the modelling of the pre-plasma is included, this ratio is reduced from ≈ 1 to a value in the range 0.7 to 0.9. Based on knowledge of the planned experimental setup, specifically that target diameters of $d = 6$ and $10 \mu\text{m}$ would be available, and that a focal spot size of $3 \mu\text{m}$ would be formed at best focus, it was clear that this optimal ratio would not be achieved under these conditions. This lead to a potential approach in which the laser is defocused to produce a larger focal spot size at the point of interaction, thereby achieving the ideal ratio, albeit at the expense of laser intensity.

The influence of spatial jitter was then investigated. Two sets of combined FLASH and EPOCH simulations were performed to understand the impact of spatial jitter magnitude on both the $2\omega_L$ yield and spatial mode. The range of magnitudes was determined by spatial jitter measurements taken from the Gemini system. These simulations demonstrate that spatial jitter can have a substantial impact on both the $2\omega_L$ yield and the generated mode shape. This effect is dependent on the focal spot size to aperture size ratio. For a $d = 6 \mu\text{m}$ aperture target, smaller magnitude spatial jitter can enhance the yield by allowing a higher intensity portion of the pulse to interact with an aperture edge. For the $d = 10 \mu\text{m}$ targets, a larger spatial jitter magnitude is required for the same effect to occur. A highly asymmetric interaction between the pulse and aperture results in the degradation of the mode quality for both aperture diameters investigated. Two potential strategies were identified to mitigate the negative impact of the spatial jitter. The first is statistical in nature, involving taking a large number of shots to increase the likelihood of an interaction occurring with the desired spatial jitter magnitude. The second involves positioning the target such that it is not at the minimum beam waist, effectively creating a ‘defocused’ condition. This would reduce the relative impact of the spatial jitter while also helping to achieve the optimal focal spot size to aperture diameter ratio. However, this approach is accompanied with a decrease in pulse intensity at the point of interaction, reducing the $2\omega_L$ yield.

Chapter 6

Experimental investigation of the generation of frequency-doubled light from pre-formed aperture targets

6.1 Introduction

This chapter reports on the results of an experimental campaign carried out to investigate the generation of frequency-doubled light from pre-formed aperture targets. This builds upon the simulation-based investigation of the generation of this form of light that is the focus of the previous two chapters of this thesis. This practical demonstration of the use of aperture targets for the generation of $2\omega_L$ light is an essential step towards its use in any potential scientific and societal applications.

The experiment was carried out in February 2022 using the Gemini laser system. It involved measuring frequency-doubled light produced by the interaction of high-intensity laser pulses ($\sim 1 \times 10^{21}$ W/cm²) with pre-formed aperture targets. The primary objective was to understand how the efficiency of frequency-doubled light generation varies with changes in aperture diameter, pulse intensity, and focal spot size. This was achieved through a comparative analysis of experimental measurements and numerical simulation results.

6.2 Experimental method

The Gemini Laser system (section 3.3.1) was used to provide pulses with a wavelength of $\lambda_L = 800$ nm, durations of $t_L \sim 45$ fs and on-target energies of up to $E_L \sim 14$ J. These energies, combined with the use of an F/2 off-axis parabola (OAP) used to bring the pulses to focus onto the aperture targets with a FWHM of ~ 3.5 μm , provided peak intensities of up to 5×10^{21} W/cm². After the F/2 focusing OAP an identical OAP was placed such that the expanded beam was collimated and directed onto a scatter screen to allow for imaging of the transmitted and generated light. A single plasma mirror (PM) (section 3.2.2) was employed prior to the main focusing optic to enhance the temporal intensity contrast by up to two orders of magnitude. This leads to a contrast of up to 10^{12} on the nanosecond scale and 10^7 on the picosecond scale prior to the arrival of the main pulse. A schematic of this setup and components is shown in figure 6.1(a).

As covered in the previous chapter, it could be advantageous to increase the FWHM of the focal spot as it interacts with an aperture target to mitigate the negative effects of any potential spatial jitter. A possible way of achieving this would be through changing the OAPs being used to those with a different F-number, but making such a large change to the beam-line during an experiment would not have been practical. Instead, the aperture targets were mounted on a drive system that allowed for movement toward the focusing OAP. Moving the aperture targets out of focus, into the incoming beam resulted in a FWHM larger than that present at best-focus when the interaction occurred. Throughout the remainder of this chapter the focal position relative to best-focus will be denoted by Δ_X , and is labelled in figure 6.1(a).

The targets employed throughout the experiment had aperture diameters of $d = 6$ or 10 μm and thicknesses of $l = 6$ μm . To create the aperture targets, the Target Fabrication Group based at the CLF developed a technique that involved etching the apertures into silicon wafers. The smallest aperture diameter available for use throughout the experiment was a consequence of the minimum diameter achievable using this technique. To aid in increasing the total number of shots taken during the experiment, multiple apertures were etched into a single silicon wafer. This wafer was then mounted onto a drive system that was used to control the position of targets to align them with the laser pulse. Once placed in the chamber, this allowed for multiple interactions to

take place without the need to perform the time consuming process of bringing the chamber from vacuum to atmospheric pressure after a single shot which would have been the case had targets with only a single aperture been used.

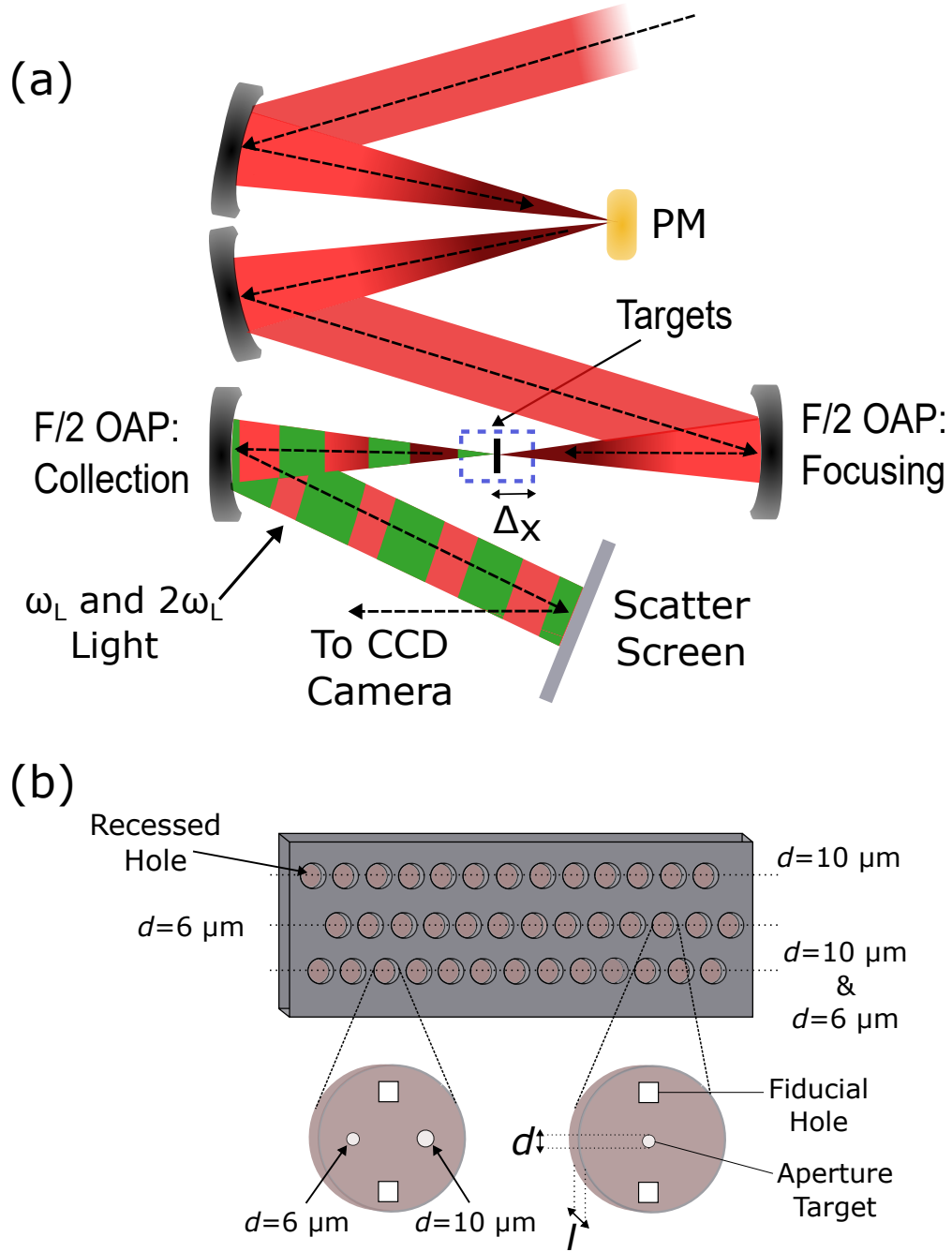


Figure 6.1: (a) Illustration of the beam-line used during the experiment. The Gemini laser beam is focused onto a plasma mirror after which the expanding beam is collimated by an identical optic. The beam is then directed onto an F/2 OAP that brings the beam to focus onto the aperture targets. An identical F/2 OAP collimates the expanding beam and directs it towards a scatter screen for imaging. (b) Illustration of a target wafer used during the experiment. Three rows of 13 recessed holes are present, within which are located the aperture targets and fiducial holes to aid alignment. The top row contains apertures of $d = 6 \mu\text{m}$, the middle row contained apertures of $d = 10 \mu\text{m}$ and the bottom row contained both a $d = 6 \mu\text{m}$ and $d = 10 \mu\text{m}$ aperture in each recessed hole.

A target wafer, a schematic of which is shown in figure 6.1(b), consisted of three rows of 13 recessed holes, within which the aperture was located. The top row contains apertures of $d = 6 \mu\text{m}$, the middle row contained apertures of $d = 10 \mu\text{m}$ and the bottom row contained both a $d = 6 \mu\text{m}$ and $d = 10 \mu\text{m}$ aperture in each recessed hole. This allowed for a maximum total of 52 shots to be taken per target wafer. In addition to the aperture targets themselves were so-called fiducial holes placed $500 \mu\text{m}$ above and below the apertures. With dimensions of $50 \times 50 \mu\text{m}$, these were included as a way of aiding in locating the apertures during the process of their alignment to the laser, prior to a shot being taken. Carried out using a highly magnified (up to $50\times$) reference camera, this process involves viewing the focal spot of the laser, in a low power mode, and manoeuvring the target such that the two are aligned. Given that the focal spot FWHM and the aperture diameter were both on the scale of microns, and the bulk of the target on the scale of centimetres, this could be a time consuming process. Hence, the fiducial holes, being on a larger scale, would be easier to locate. Once located, finding the aperture target was a short process due to the distance between it and the fiducial being known. This prevented the alignment process from becoming a limiting factor on the total number of shots taken throughout the experiment. The fiducial holes were placed both above and below the apertures due to the possibility that a different diameter target would be used from one shot to the next, requiring a move to a target in a row above or below that which was just used.

6.2.1 Diagnostics

To obtain a measure of the total photon count of the light both generated from, and transmitted through, the aperture targets, a diffuse scatter screen composed of PTFE was placed 40 cm from the collection OAP such that the fully expanded beam was incident upon it, shown in figure 6.1(a). Two CCD cameras fitted with $1.2\times$ telephoto lenses were then used to image the light on the scatter screen, one filtered for ω_L light, the other for $2\omega_L$ light. In order to determine the correct level of ND filtering required for the cameras, a series of shots without targets present at 50% of the maximum energy were performed. ND filtering was then applied such that the maximum signal measured on the ω_L CCD camera was in the center of the dynamic range of the camera. This energy was chosen with the assumption that in the case of the pulse propagating through the center of the aperture target, some of the pulse energy would be blocked or

absorbed by the bulk of the target. A similar level of ND filtering was initially applied to the $2\omega_L$ camera. A high energy $2\omega_L$ laser was not available for use in a no-target shot, therefore the ND level would later be adjusted if required upon the measurement of $2\omega_L$ light from initial target shots.

6.3 Image masking

To measure the detected light on the scatter screen, the total signal present in each captured image was evaluated. As the scatter screens for these diagnostics had to be positioned in such a way that the transmitted and generated light would be incident upon them, and due to the limited number of windows available for viewing into the chamber, the CCD cameras also captured spurious light from inside the chamber. Present in the majority of the images are reflections of light from other objects, such as metal optic mounts or the chamber walls, that could potentially contribute to the measured signal.

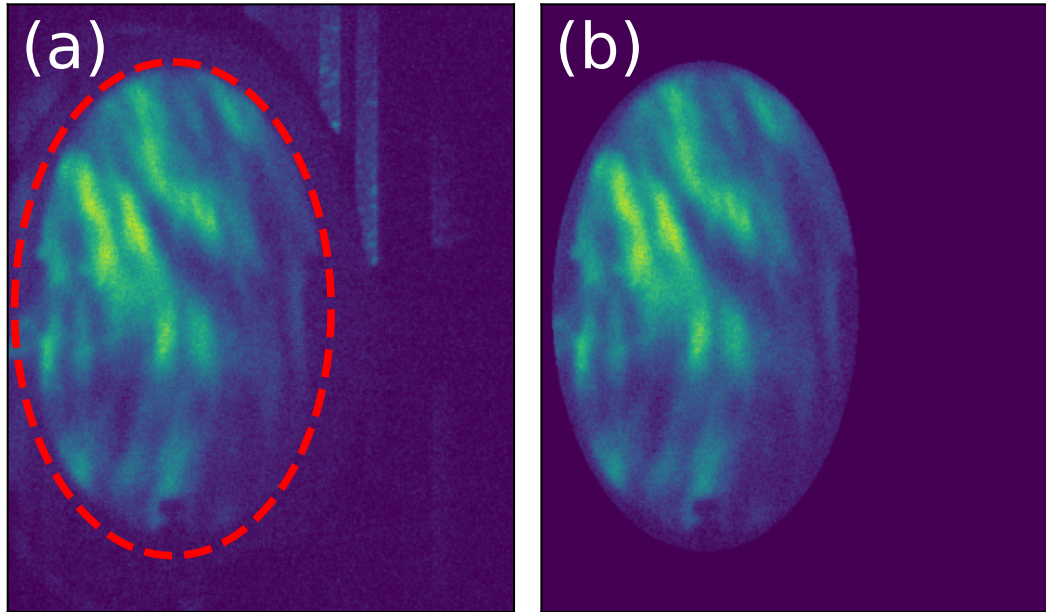


Figure 6.2: Examples of an image captured by the transmission ω_L diagnostic (a) before and (b) after the masking process used to remove unwanted signal from the image.

To address this, a mask was applied to the images that selected only the portion of the image containing the scatter screen. This then ensured that only the light from the interaction that was incident on the scatter screen is present when the signal is summed. An example of an image before and after masking is shown in figures 6.2(a) and (b) respectively. Note the background reflections in the top right of (a) that have been removed by the masking process. Minor movements of the camera or scatter screen throughout the experiment lead to slight changes in the exact position of the light as it was incident on the scatter screen. As such, the mask position was manually adjusted on an image-by-image basis, as required.

6.4 Focal spot image analysis

As mentioned above, the focal spot diameter was varied throughout the experiment via altering the Δ_X position of the target as a way of potentially mitigating the impact of the on-shot spatial jitter magnitude on $2\omega_L$ light generation, as discussed in the previous chapter. Three values were chosen; $\Delta_X = 0 \mu\text{m}$, $20 \mu\text{m}$ and $29 \mu\text{m}$, with the target moved toward the focusing optic in the latter two cases, as indicated in figure 6.1(a). To understand the results obtained during the experiment, it is important to characterise the focal spot spatial-intensity distributions for each of the selected Δ_X values. This was enabled by capturing images of the focal spot spatial-intensity distributions during the alignment process mentioned above, using the reference camera.

Figure 6.3 shows example focal spot intensity distributions measured for each Δ_X position. In the case of $\Delta_X = 0 \mu\text{m}$, a focal spot with an approximately Gaussian distribution was observed, with some variation in the quantity of light present around the edges. However, as Δ_X increases, the focal spot shape is seen to become increasingly non-Gaussian, with the spatial distributions measured for the cases of $\Delta_X = 20 \mu\text{m}$ and $29 \mu\text{m}$ showing significantly more light present outside the central intensity region. While these examples are representative, focal spot intensity distributions measured for the same defocus position did display some variation. Had focal spot images been capture prior to every shot taken, there would have been the potential to directly correlate the size and shape with the experimental measurements on a shot-by-shot basis. However, due to the additional time required to capture these images, the decision was made to prioritise taking a greater number of shots.

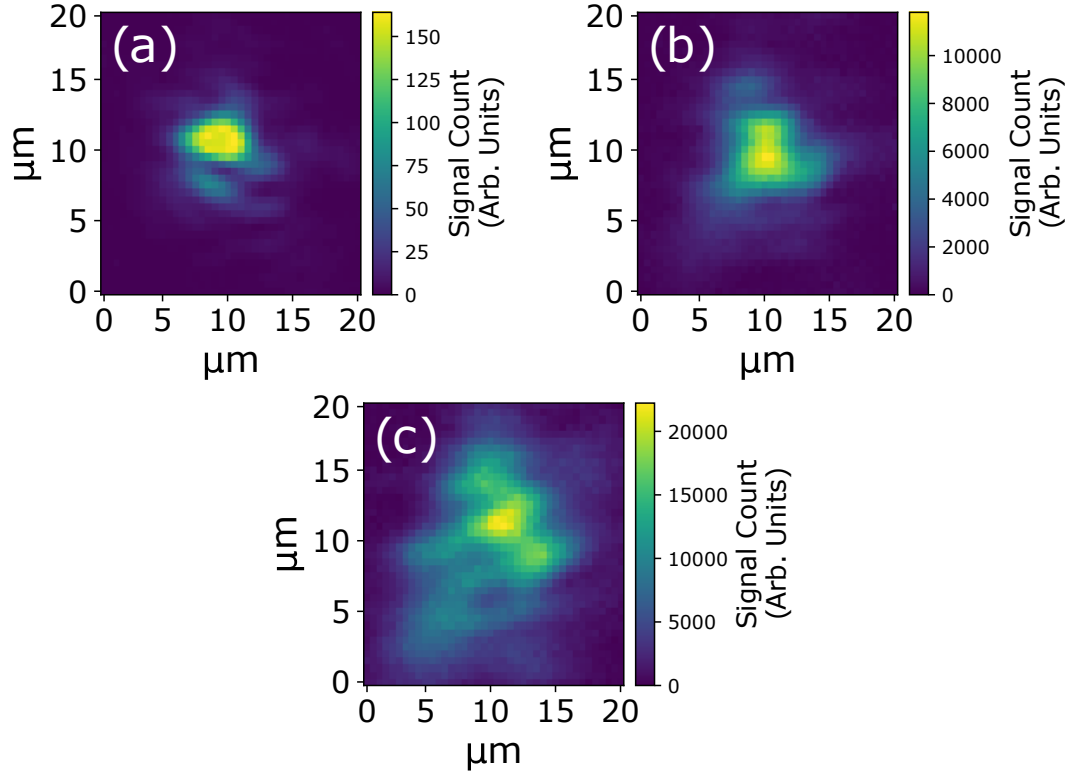


Figure 6.3: Example focal spot measurements obtained during the experiment for Δ_X values of (a) 0 μm ; (b) 20 μm and (c) 29 μm .

6.5 Simulations performed to aid in understanding the experimental results

To aid in understanding the measurements made during the experiment, combined FLASH and EPOCH simulations were performed. The FLASH simulations were performed under the same conditions discussed in the previous chapter. For the EPOCH simulations, the grid was defined as 3200×3200 mesh cells, corresponding to a simulation box size of $45 \mu\text{m} \times 45 \mu\text{m}$, with a particle-per-cell number of 15. A pre-plasma profile generated from a FLASH simulation was directly imported and used to define the ion and electron density distributions. A linearly polarised pulse with a wavelength of $\lambda_L = 800 \text{ nm}$ was used, entering the simulation box from the left (X) boundary, with a duration of $t_L = 40 \text{ fs}$ (FWHM). The simulations were performed for 200 fs. All boundaries were configured as free space. The focal spot size at the point of interaction was varied within the simulations by changing the target Δ_X value, with $\Delta_X = 0$,

20 and 29 μm used. The pulse intensity was varied between $1.75 \times 10^{20} \text{ W/cm}^2$ and $5 \times 10^{21} \text{ W/cm}^2$ by increasing the pulse energy. Apertures with diameters of $d = 6 \mu\text{m}$ and $10 \mu\text{m}$ were used and with a target thickness of $l = 6 \mu\text{m}$. The position of the focal spot was varied between the center of the aperture to a maximum distance of $9.5 \mu\text{m}$ to mimic the effect of the spatial jitter on the interaction.

6.6 Experimental results

This section explores the measurements made during the experiment, investigating the way in which $2\omega_L$ varies as a function of on-target energy, ω_L signal, aperture diameter and target Δ_X position, understood through comparison with the results from the simulations detailed above.

6.6.1 Variation of detected ω_L and $2\omega_L$ light as a function of on-target energy

The simulation-based investigation of the generation of $2\omega_L$ light discussed in the previous two chapters demonstrate that a difference in the measured $2\omega_L$ signal would be expected as the interaction parameters of pulse energy and focal spot size are varied. Figures 6.4(a) and (b) show the variation in measured $2\omega_L$ signal with on-target energy, E_L , for targets of $d = 6$ and $10 \mu\text{m}$ for $\Delta_X = 0$ and $29 \mu\text{m}$ respectively. The uncertainties in the pulse energy are quoted from the Gemini laser facility and those in the summed signal arise from calibrating the light level on the transmission screen used to make the measurements.

As can be seen, the measured $2\omega_L$ signals show no discernable trend. This suggests that interaction parameters other than pulse energy and focal spot size play a dominant role in the generation of frequency doubled light. To investigate this further, the measurements made for $\Delta_X = 0$ case were compared with simulation data.

Figures 6.5(a) and (b) show the changes in the measured $2\omega_L$ signal as a function of on-target energy for $d = 6 \mu\text{m}$ and $d = 10 \mu\text{m}$ respectively, both for $\Delta_X = 0 \mu\text{m}$. Plotted with these measurements are the $2\omega_L$ signals from 2D simulations in which E_L was varied for otherwise fixed target and pulse parameters, corresponding to the experimental parameters. In all simulations within this section, the greatest maximum pulse intensity that was used was $5 \times 10^{21} \text{ W/cm}^2$, chosen to match the maximum

pulse intensity achieved during the experiment. The values on both axes have been normalised to enable direct comparison. A limitation of this approach is that only the experimental Δ_X values were used. As 2D simulations were performed, the intensity does not reduce with Δ_X as rapidly as occurred during the experiment.

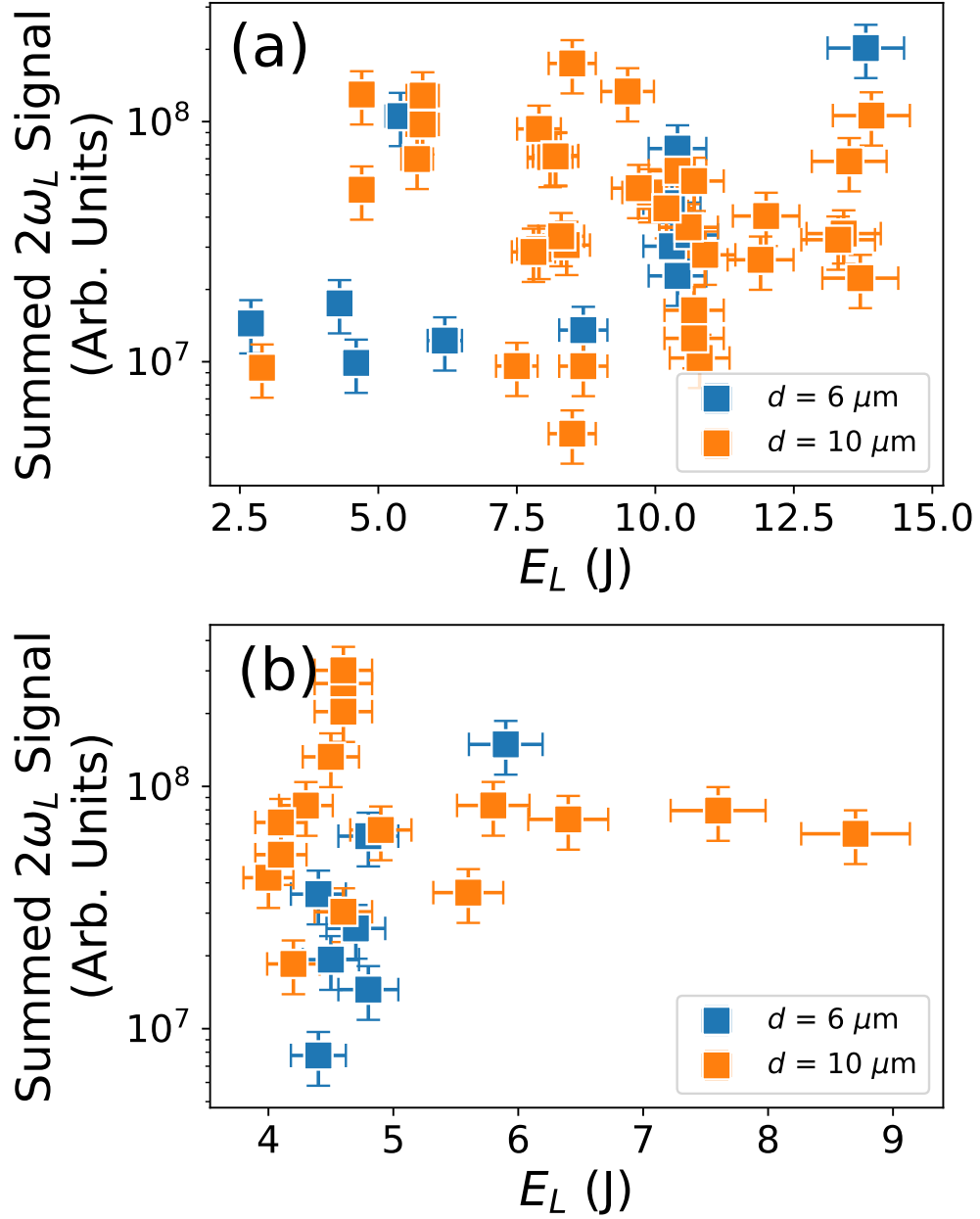


Figure 6.4: $2\omega_L$ signal as measured during the experiment for $d = 6$ and $10 \mu\text{m}$ for Δ_X values of (a) $0 \mu\text{m}$ and (b) $29 \mu\text{m}$.

From the simulation results in figures 6.5(a) and (b), an increase in the $2\omega_L$ signal with increasing pulse energy can be observed. This can be attributed to the increased energy available to accelerate electron bunches of greater density from the aperture edges, generating greater quantities of frequency-doubled light. The same trend is not observed in the experimental measurements. While figure 6.5(a) shows a general trend of increasing $2\omega_L$ signal with on-target energy, there are clear outlying measurements, notably the cluster for which the same energy, $E_{L_{Norm}} \sim 0.8$, results in a $2\omega_L$ signal that varies by $3.3\times$ from the minimum to maximum value. Next, considering the range of experimentally measured $2\omega_L$ values displayed in this figure, a change in the signal of roughly $20\times$ from its minimum to maximum value is seen. The on-target energies generating these $2\omega_L$ signals ranged between 2.7 and 13.8 J, representing a roughly $5\times$ increase in pulse energy. To capture a similar change in $2\omega_L$ from the simulations, a much greater range of pulse intensities was required, with this parameter being varied from 2.5×10^{20} W/cm² to 5×10^{21} W/cm², representing a $20\times$ change. As these are 2D simulations, the reported intensities do not directly correspond to experimental values without a dimensionality correction. However, this and other similar comparisons made in this chapter focus on the relative variation required to reproduce the observed range of $2\omega_L$ signals. As such, this correction would not affect the conclusion that a significantly larger range of pulse intensities is required in simulation than in experiment.

The $2\omega_L$ experimental measurements shown in figure 6.5(b), for targets with $d = 10$ μm , appear more randomly scattered, not showing an increase with pulse energy as seen from the simulation results. Similarly to figure 6.5(a) are clusters of points for which a similar pulse energy, $E_{L_{Norm}} \sim 0.8$, produces a range of $2\omega_L$ signals, an approximately $6\times$ increase from the minimum to maximum in this case. Again considering the range of $2\omega_L$ signal measured, a $34\times$ increase was induced by changes in E_L from 2.9 J to 13.9 J, a roughly $4.7\times$ increase in energy. Generating a similar range of $2\omega_L$ values within the simulations required the pulse intensity to be varied from 1.75×10^{20} W/cm² to 5×10^{21} W/cm², a approximately $28\times$ increase from minimum to maximum. There is a clear disparity between the way in which the $2\omega_L$ signal varies with E_L when comparing the experimental measurements and the simulation results. In addition, there are notably different ranges of pulse energies required to generate a similar range of $2\omega_L$ signals. These observations strongly suggests that factors other than pulse energy played a

dominant role in the quantity of $2\omega_L$ light generated from an interaction during the experiment.

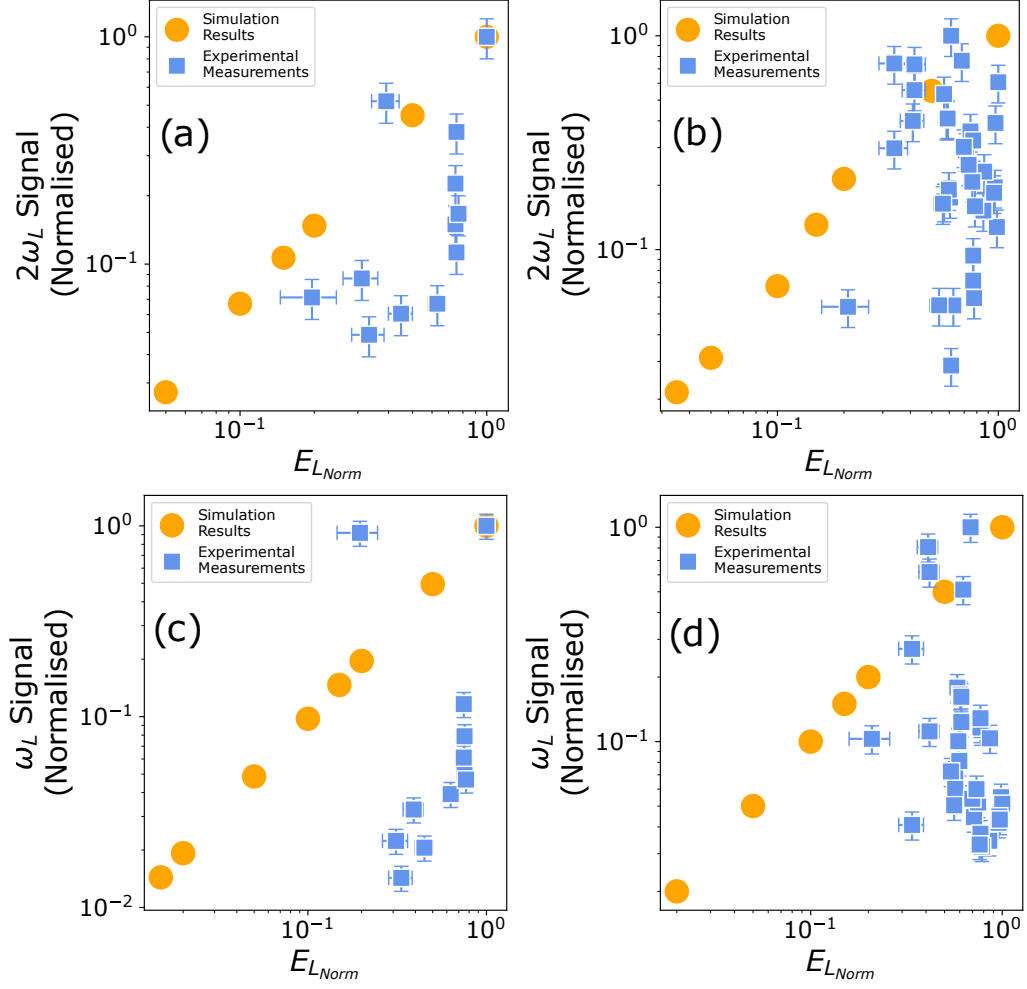


Figure 6.5: Measured and simulated signals as a function of on-target energy: the $2\omega_L$ signal for (a) $d = 6 \mu\text{m}$ and (b) $d = 10 \mu\text{m}$, and the ω_L signal for (c) $d = 6 \mu\text{m}$ and (d) $d = 10 \mu\text{m}$. All cases use additional target parameters $\Delta_X = 0 \mu\text{m}$ and $l = 6 \mu\text{m}$.

Given previous simulation-based modelling of the interaction presented in the previous chapter, the on-shot spatial jitter could give rise to significant variation in $2\omega_L$ emission. If this were the case, variations due to the spatial jitter would also be observed in the ω_L transmitted light as it varies with pulse energy, with these variations attributable to the changes in the percentage of the pulse that is blocked by the bulk of the target as the spatial jitter magnitude changes. These measurements, along with accompanying simulation results, are shown in figures 6.5(c) and (d) for $\Delta_X = 6 \mu\text{m}$ and $10 \mu\text{m}$, respectively.

The simulation results show an increase in ω_L signal with increasing energy as would be expected. The experimental measurements of ω_L shown in figure 6.5(c) display a general increase with on-target energy, similar to that seen for the $2\omega_L$ signal, with similar outlying clusters of measurements such as a range of signals resulting from similar on-target energies. Again, the range of energies required to generate similar ω_L signal differs greatly between simulation and experiment, with $66\times$ and $5\times$ required, respectively. In figure 6.5(d), the experimental measurements seem randomly distributed, similar to the $2\omega_L$ measurements made under the same interaction conditions, clearly not in agreement with the simulation results. Once again, the on-target energy required to produce a similar $2\omega_L$ range differed markedly: the experimental measurement a $\sim 4.7\times$ increase while the simulations required $\sim 50\times$ increase.

The analysis of the ω_L measurements supports the hypothesis that spatial jitter magnitude may be a factor in explaining the change in the experimental measurements of the $2\omega_L$ signal seen in figure 6.5. Plots in which ω_L and $2\omega_L$ are shown as a function of on-target energy for the remaining two Δ_X positions reveal a similar lack of agreement between the simulation results and experimental measurements, demonstrating that this is the case for the measurements made under all interaction conditions investigated. Based on these findings, the decision was made to normalise all measured signals to pulse energy in the subsequent analysis, denoted by $2\omega_{LNorm}$. This allowed for the effects of other variables, such as spatial jitter magnitude, to be investigated.

6.6.2 Fundamental and second harmonic signal correlation with spatial jitter magnitude

With the previous section giving evidence that the on-shot spatial jitter magnitude was likely an important factor in the generation of frequency-doubled light, the next stage was to investigate the $2\omega_L$ signal as a function of ω_L signal to determine if the trends observed could be understood by considering its effect on the interaction. As mentioned in the previous chapter, the spatial jitter magnitude was measured for the Gemini system, using a low-power mode, prior to the experimental campaign. Figure 6.6 shows the histogram of these measurements, demonstrating that an average magnitude of $1.6\ \mu\text{m}$ and a maximum magnitude of $5\ \mu\text{m}$ were observed. Throughout this section, in the simulations performed to understand the experimental measurements, the spatial jitter magnitude was extended such that a similar range of $2\omega_L$ signals was observed,

often resulting in values greater than $5\ \mu\text{m}$ being used.

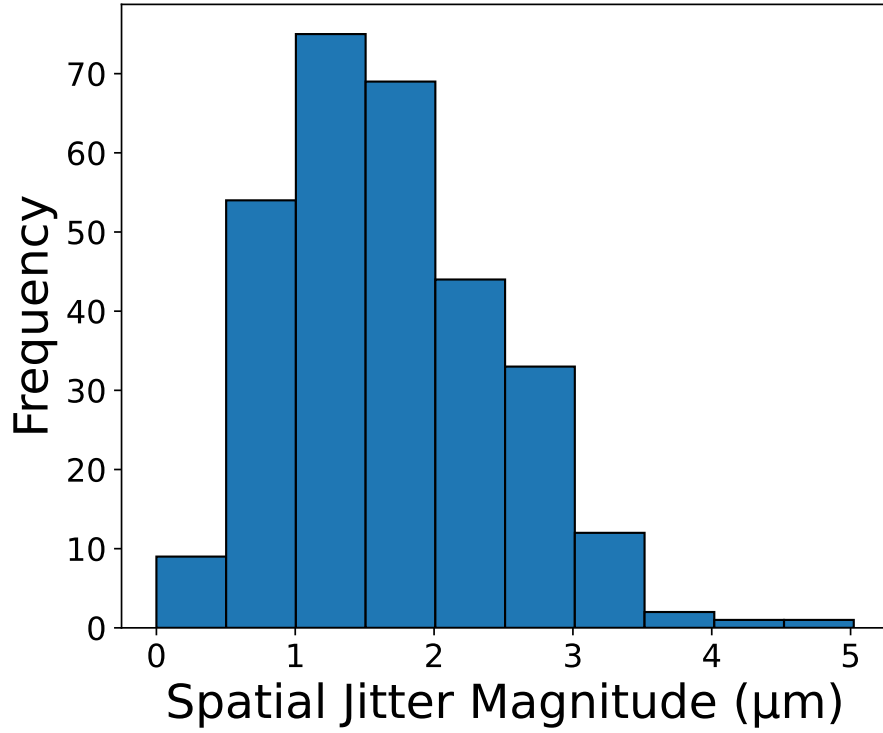


Figure 6.6: Histogram displaying the frequency of spatial jitter magnitude over 300 shots, taken using a low-power mode of the Gemini Laser System.

Results for the case of $d = 6\ \mu\text{m}$, $\Delta_X = 0\ \mu\text{m}$

Figure 6.7(a) shows the measured $2\omega_L$ signal as a function of ω_L signal, for $\Delta_X = 0\ \mu\text{m}$ and for targets with $d = 6\ \mu\text{m}$, both normalised to pulse energy. Also displayed are the results of a set of simulations performed with the same target and pulse parameters in which the spatial jitter magnitude has been varied from 0 to $4.5\ \mu\text{m}$. Both experimental measurements and simulation results have then been normalised to enable direct comparison.

Observed in the simulation results is an increase in $2\omega_L$ generation with increasing spatial jitter magnitude until a maximum is reached at a value of $1.5\ \mu\text{m}$. As discussed in the previous chapter, this can be explained by fact that a higher intensity portion of the pulse interacts with the aperture edge as the focal spot of the laser pulse moves from the center of the aperture toward one of the edges.

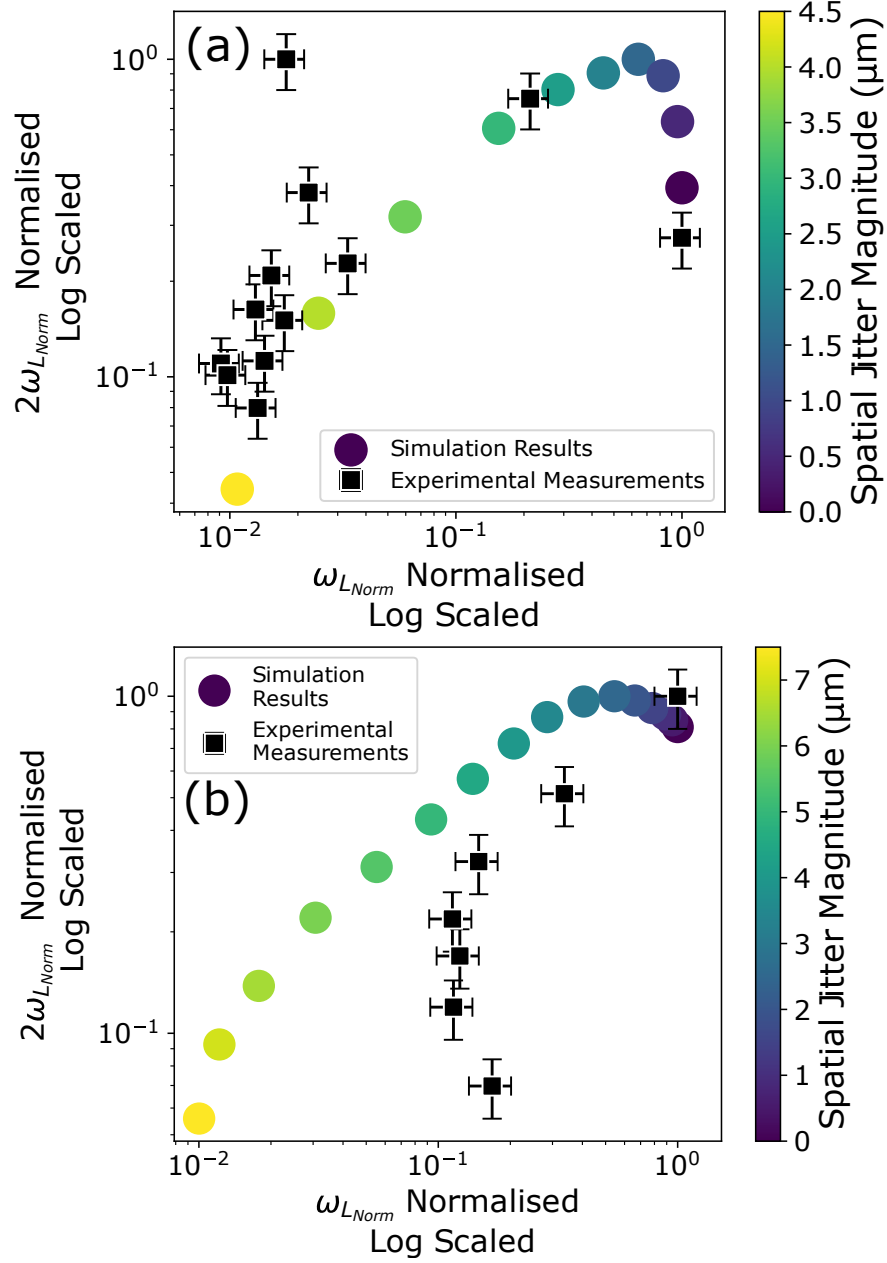


Figure 6.7: $2\omega_L$ experimental measurements as a function of ω_L signal for targets with $d = 6 \mu\text{m}$ for Δ_x values of (a) $0 \mu\text{m}$ and (b) $29 \mu\text{m}$. Also plotted are the same metrics observed in combined FLASH and EPOCH simulations performed in which the spatial jitter of the laser was varied from (a) $0 \mu\text{m}$ to $4.5 \mu\text{m}$ and (b) $0 \mu\text{m}$ to $7.5 \mu\text{m}$ for the same respective target conditions as for the experimental measurements.

The increase in intensity at the point of interaction leads to the generation of higher density electron bunches and therefore a greater quantity of frequency-doubled light. This is accompanied by a slight decrease in the quantity of ω_L light observed as a greater portion of the pulse is blocked by the bulk of the target. As the spatial jitter magnitude

increases beyond the value of $1.5 \mu\text{m}$, both the ω_L and $2\omega_L$ light observed begins to decrease. This can be attributed to the fact that the center, which is also the highest intensity portion, of the pulse moves past the aperture edge such that it interacts with the bulk of the target. This reduces the intensity at the point of interaction with the aperture edge, decreasing $2\omega_L$ generation, and blocks a larger portion of the pulse, thus reducing the observed ω_L light.

As can be seen, there is some agreement between the experimental measurements and the simulation data. While there are clear outliers, notably the measurement for which the highest $2\omega_L$ signal was observed, the other experimental measurements seem to follow a similar trend to that seen in the simulation results. This figure also suggests that a majority of the shots taken under these interaction conditions occurred at a spatial jitter magnitude of $> \sim 2.5 \mu\text{m}$, as demonstrated by the cluster of measurements that sit near the tail of the simulation results at this spatial jitter magnitude. As the simulations presented here only include a single variable, spatial jitter magnitude, the scatter of the points suggest that there are other factors influencing the generation of $2\omega_L$.

Results for the case of $d = 6 \mu\text{m}$, $\Delta_X = 29 \mu\text{m}$

Figure 6.7(b) shows the measured $2\omega_L$ signal as a function of ω_L signal for $\Delta_X = 29 \mu\text{m}$ and for targets with $d = 6 \mu\text{m}$ and $l = 6 \mu\text{m}$, both normalised to pulse energy, again displayed with the results of a set of simulations performed with the same target and pulse parameters. In this set of simulations, the spatial jitter magnitude has been varied from 0 to $7.5 \mu\text{m}$. Both experimental measurements and simulation results have then been normalised to enable direct comparison.

A similar trend can be observed in the simulation results as seen in figure 6.7(b), with an increase in $2\omega_L$ with spatial jitter magnitude before reaching a maximum value, then decreasing with further movement of the focal spot into the bulk of the target. The focal spot of the pulse at this Δ_X value is larger and broader, shown in figure 6.8, leading to the much more gradual increase seen in $2\omega_L$ signal with spatial jitter magnitude when compared with figure 6.7(a). This can be attributed to the fact that the change in intensity of the pulse at the point of interaction between the pulse and aperture edge as a function of spatial jitter magnitude is much more gradual in this case. While the $2\omega_L$ signal reaches its maximum value for a slightly greater spatial

jitter magnitude of $2.5 \mu\text{m}$ than for the case of $\Delta_X = 0$, the difference in $2\omega_L$ signal at this spatial jitter magnitude and this value at the surrounding points is minimal. An additional consequence of the larger focal spot size is the fact that, to capture a similar range of $2\omega_L$ values, a greater range of spatial jitter magnitudes are required, extending to $7.5 \mu\text{m}$ in this case.

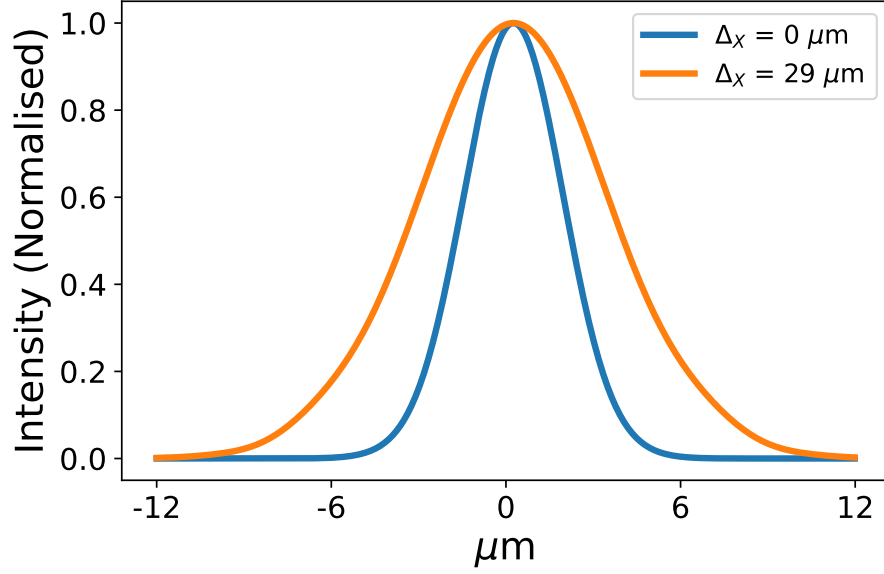


Figure 6.8: Line-outs taken through the focal spots of pulses from 2D EPOCH simulations for $\Delta_X = 0 \mu\text{m}$ and $29 \mu\text{m}$.

The experimental measurements displayed in figure 6.7(b) show a faster drop in $2\omega_L$ with decreasing ω_L light than observed in the simulation results. This lack of correlation suggests that factors other than spatial jitter magnitude that are influencing the $2\omega_L$ generation are playing a significant role. As discussed in the previous section, the focal spot spatial intensity profile is notably less Gaussian than in the case of $\Delta_X = 0$, so this likely plays a significant role in the interaction. As this is a factor that was not included in the simulations it may be the cause of the difference between simulation results and experimental measurements. Additionally, for this value of Δ_X , only 7 repeat measurements were made and, due the statistical nature of the spatial jitter, more data may be required to fully explore this behaviour.

Results for the case of $d = 10 \mu\text{m}$, $\Delta_X = 0 \mu\text{m}$

Figure 6.9(a) shows the measured $2\omega_L$ signal as a function of ω_L signal for $\Delta_X = 0 \mu\text{m}$ and for targets with $d = 10 \mu\text{m}$, both normalised to pulse energy. Also displayed are the

results of a set of simulations performed with the same target and pulse parameters in which the spatial jitter magnitude has been varied from 0 to 7.5 μm . Both experimental measurements and simulation results have then been normalised to enable direct comparison.

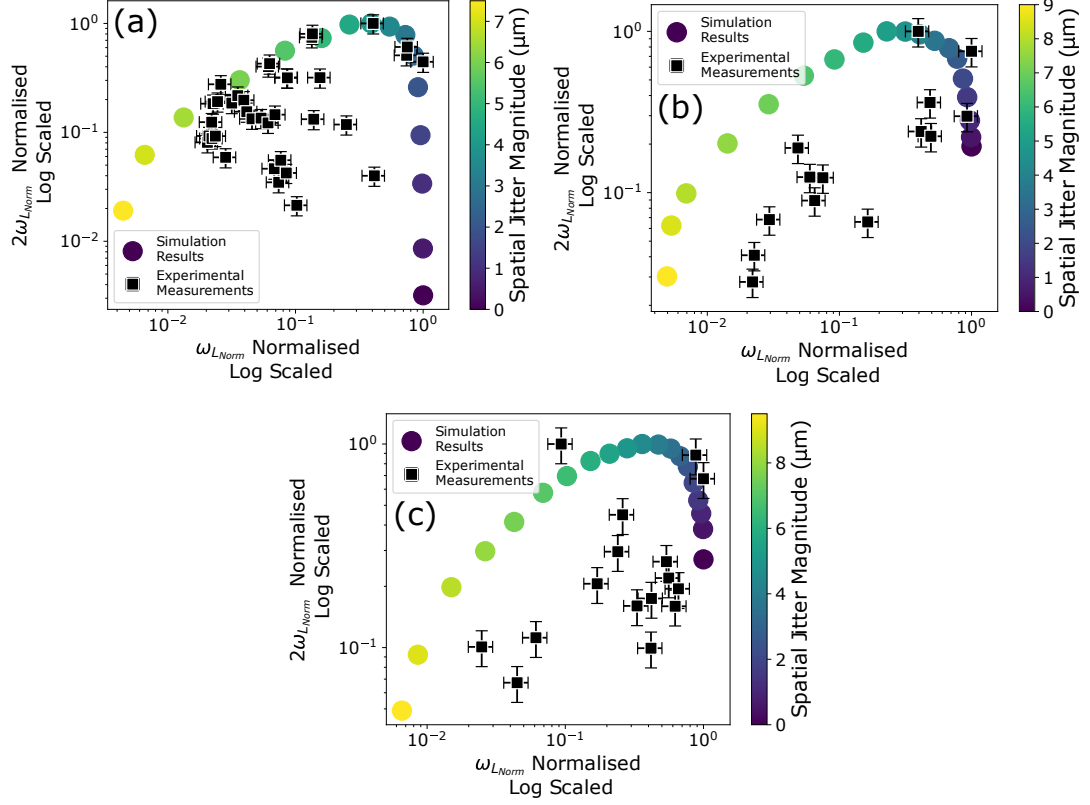


Figure 6.9: $2\omega_L$ experimental measurements as a function of ω_L signal for targets with $d = 6 \mu\text{m}$ for Δ_X values of (a) 0 μm , (b) 20 μm and (c) 29 μm . Also plotted are the same metrics observed in combined FLASH and EPOCH simulations performed in which the spatial jitter of the laser was varied from (a) 0 μm to 7.5 μm , (b) 0 μm to 9 μm and (c) 0 μm to 9.5 μm for the same respective target conditions as for the experimental measurements

Again, the simulation results show a similar trend of reaching a maximum $2\omega_L$ signal, in this case at a spatial jitter magnitude of 4.5 μm , after which a decrease in both $2\omega_L$ and ω_L signal is observed. The spatial jitter magnitude at which the maximum $2\omega_L$ signal is observed is larger than for the use of targets with $d = 6 \mu\text{m}$ due to the greater distance that the focal spot must travel before its point of maximum intensity interacts with the aperture edge.

With respect to the experimental measurements, there is some agreement between these and the simulation results in that there is a series of these values that align with the simulation results. Somewhat similar to the case of $d = 6 \mu\text{m}$ and $\Delta_X = 0 \mu\text{m}$,

there is a cluster of measurements that sit near the simulation results at a greater spatial jitter magnitude but in this case, the magnitude is greater, around $\sim 6 \mu\text{m}$. However, there is a great deal of scatter amongst the measurements that lie beneath the simulation data, given that these measurements were made for $\Delta_X = 0 \mu\text{m}$, this is unlikely to be due to the focal spot quality. It may be attributable to other factors previously mentioned such as variations in pulse contrast. In addition, there are a greater number of measurements made under these conditions when compared with those in figure 6.7(a), possibly allowing for the ‘noise’ in the signal to become more apparent.

Results for the case of $d = 10 \mu\text{m}$, $\Delta_X = 20 \mu\text{m}$

Figure 6.9(b) shows the measured $2\omega_L$ signal as a function of ω_L signal for $\Delta_X = 20 \mu\text{m}$ and for targets with $d = 10 \mu\text{m}$ and $l = 6 \mu\text{m}$, both normalised to pulse energy. Also displayed are the results of a set of simulations performed with the same target and pulse parameters in which the spatial jitter magnitude has been varied from 0 to $9 \mu\text{m}$. Both experimental measurements and simulation results have then been normalised to enable direct comparison.

The simulation results are similar to those displayed in the previous figures, with the spatial jitter magnitude at which the maximum $2\omega_L$ signal is observed being at $5 \mu\text{m}$, similar to the previous figure. A more gradual increase is seen in the $2\omega_L$ signal with increasing spatial jitter magnitude, compared to that seen in figure 6.9(a). Again, this is attributable to the broader focal spot profile present at the point of interaction for this Δ_X value.

Considering the experimental measurements, there is a general increase in $2\omega_L$ signal with ω_L . However, they do not follow the same trend seen in the simulation results. This can be in part attributed to the likely non-Gaussian focal spot profile present at this Δ_X value.

Results for the case of $d = 10 \mu\text{m}$, $\Delta_X = 29 \mu\text{m}$

Figure 6.9(c) shows the measured $2\omega_L$ signal as a function of ω_L signal for $\Delta_X = 29 \mu\text{m}$ and for targets with $d = 10 \mu\text{m}$ and $l = 6 \mu\text{m}$, both normalised to pulse energy. Also displayed are the results of a set of simulations performed with the same target and pulse parameters in which the spatial jitter magnitude has been varied from 0 to $9 \mu\text{m}$.

Both experimental measurements and simulation results have then been normalised to enable direct comparison.

The experimental measurements do show a general increase in $2\omega_L$ with ω_L signal but, as with the previous figure, the trend is not the same as that seen in the simulation measurements. Similar to the case of $\Delta_X = 20 \mu\text{m}$ and in the previous plot of $\Delta_X = 29 \mu\text{m}$ for targets with $d = 6 \mu\text{m}$, this can likely be explained in part by the non-Gaussian focal spot spatial profile present at this defocus value.

Overall, the best agreement between experimental measurements and simulation results is observed for the case of $\Delta_X = 0 \mu\text{m}$ for both aperture diameters, $d = 6$ and $10 \mu\text{m}$. For these interaction conditions, there is a notable degree of scatter in the measurements. This suggests that, while spatial jitter magnitude may have been a major determining factor in the quantity of $2\omega_L$ light generated, other factors still played a significant role in the interaction dynamics. For the cases of $\Delta_X > 0$, less agreement between the simulation results and experimental measurements is observed. This can be likely be attributed to the focal spot spatial profiles departure from the Gaussian shape predicted by simple Gaussian optics and could potentially produce the scatter that is seen in the measurements.

6.6.3 Second harmonic signal as a function of aperture diameter

In this section, the experimental measurements made using the two aperture diameters are compared with the results of simulations presented in the previous section. For the following analysis, the spatial jitter magnitude has been limited to $5 \mu\text{m}$, in line with the magnitude measured prior to the beginning of the experiment, shown in figure 6.6. Sufficient numbers of measurements were made at $\Delta_X = 0 \mu\text{m}$ and $29 \mu\text{m}$ for analysis to be performed.

Results for the case of $\Delta_X = 0 \mu\text{m}$

Examining first the case of $\Delta_X = 0 \mu\text{m}$, figure 6.10(a) shows the change in $2\omega_L$ signal as a function of spatial jitter magnitude for $d = 6$ and $10 \mu\text{m}$, as a result of two sets of simulations. As can be seen, a similar maximum signal is observed for both values of d and minimum signal is produced when a $d = 10 \mu\text{m}$ target is used for this range of spatial jitter magnitudes. Figure 6.10(b) shows the experimental measurements displayed alongside the range of the simulation results as a function of d . The simulation

results, represented by the green shaded region, have been normalised to the experimental measurements for each aperture diameter to enable direct comparison. The experimental measurements are represented by a point indicating their mean value, with the red shaded region displaying the range of signals measured, and the ‘hatched’ box displaying the range of one standard deviation of the measurements around the mean. Simulation results are only single values for each spatial jitter magnitude, hence no mean or standard deviation has been included.

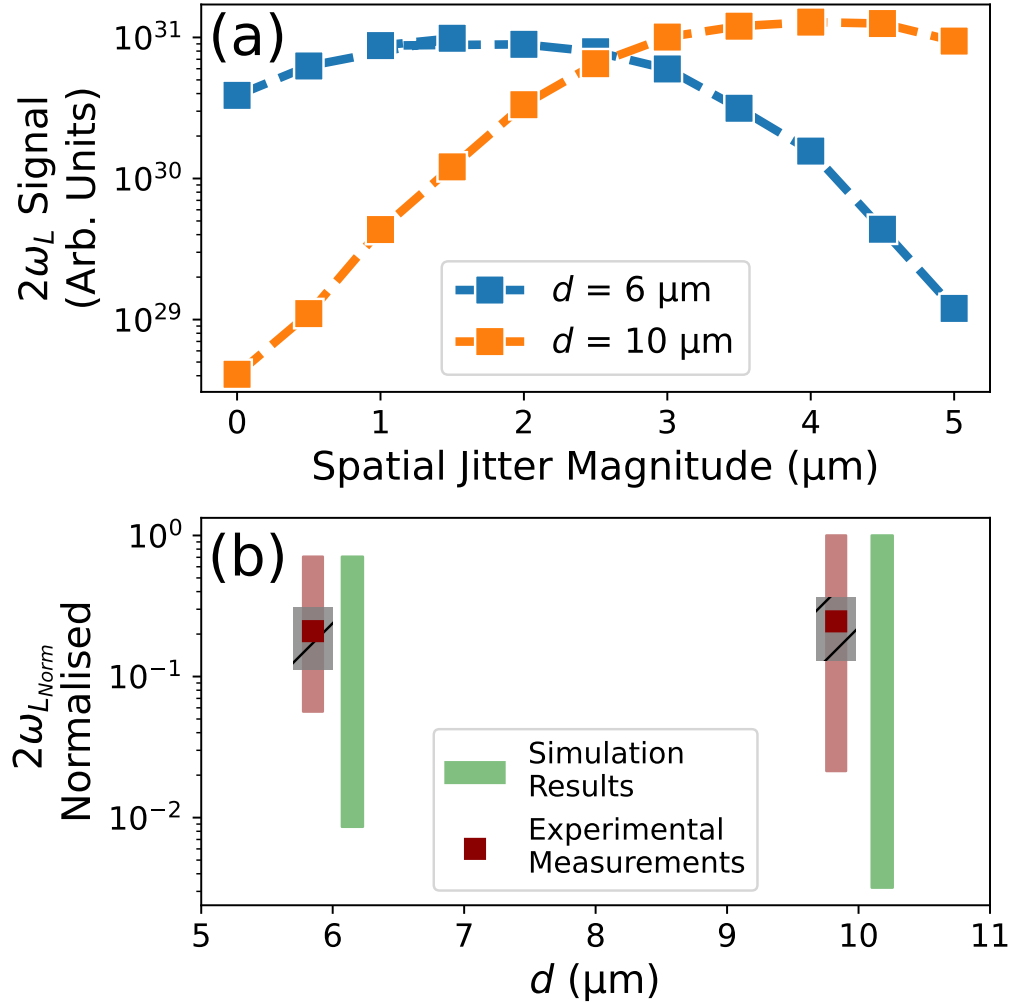


Figure 6.10: (a) Results of the two sets of simulations performed in which spatial jitter magnitude was varied for targets with $d = 6 \mu\text{m}$ and $10 \mu\text{m}$ for $\Delta_X = 0 \mu\text{m}$. (b) The experimental measurements, indicated by a point for the mean value, with the range of values represented by the red shaded region and the ‘hatched’ box representing the standard deviation, compared with the range of $2\omega_L$ signals from the simulations shown in (a), represented by the green shaded region.

Overall, agreement is seen between the ranges of both the simulation results and experimental measurements under these conditions. While the range of the simulation results clearly extends beyond the experimental measurements for both aperture diameters, both the maximum and minimum signal was measured for a $d = 10 \mu\text{m}$ target, aligning with the simulation results. The difference in the range of signal between experimental measurement and simulation results could stem from the spatial jitter magnitude being concentrated around the range of ~ 0.5 to $\sim 3 \mu\text{m}$, as seen from the histogram in figure 6.6, as this could potentially provide the measurements observed here.

Results for the case of $\Delta_X = 29 \mu\text{m}$

Now considering the case of $\Delta_X = 29 \mu\text{m}$, figure 6.11(a) shows the change in $2\omega_L$ signal as a function of spatial jitter magnitude for $d = 6$ and $10 \mu\text{m}$, as a result of two sets of simulations. As can be seen, approximately equal maximum signals are observed for both aperture diameters, and the minimum signal is observed for the use of a $d = 10 \mu\text{m}$ target for this range of spatial jitter magnitudes. Shown in figure 6.10(b) are the experimental measurements displayed alongside the range of the simulation results, as a function of d , displayed in the same manner as in the previous figure.

For this Δ_X position, the spatial jitter modelled by the simulations does not fully account for the $2\omega_L$ signal measured during the experiment. For both aperture diameters, the range of measured $2\omega_L$ signals exceeds that observed in the simulation results. As suggested in previous sections, this may be largely attributable to the non-Gaussian focal spot spatial intensity profile that exists at the point of interaction for this Δ_X value that is not modelled by simulations.

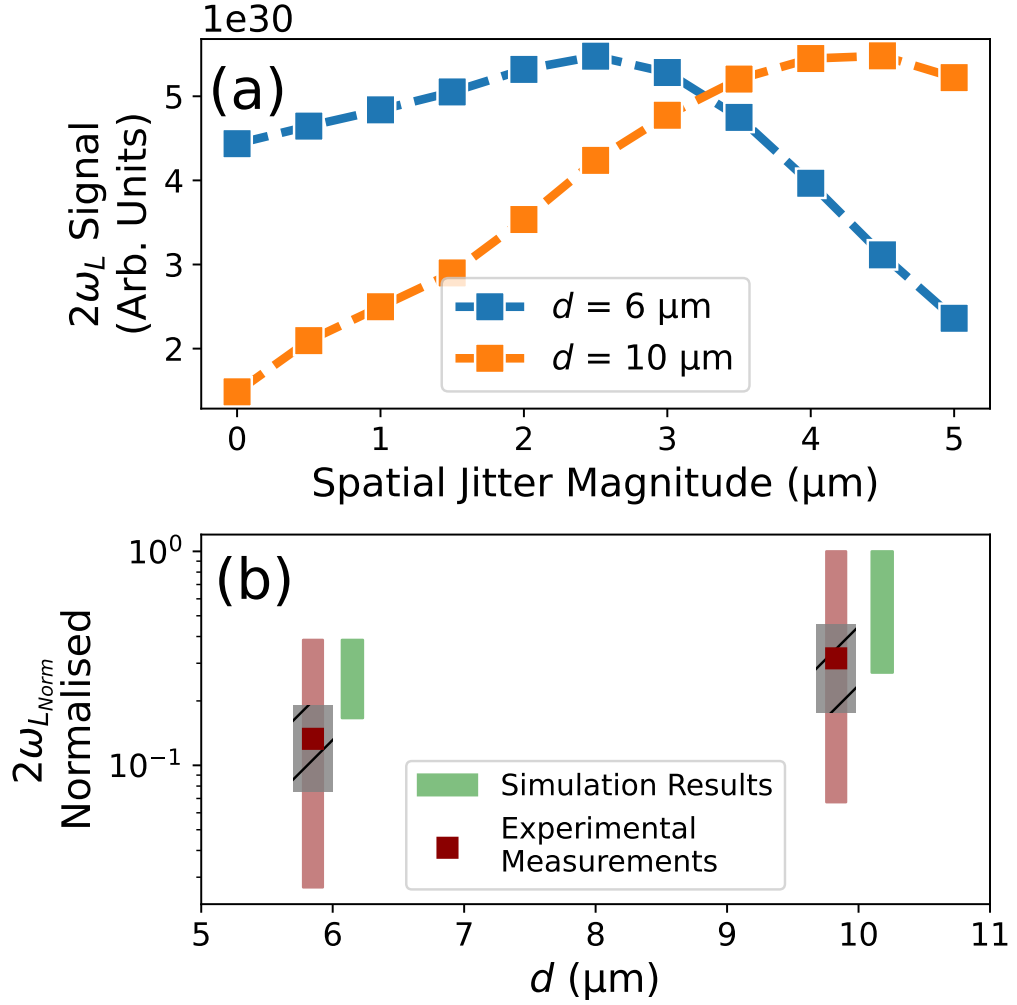


Figure 6.11: (a) Results of the two sets of simulations performed in which spatial jitter magnitude was varied for targets with $d = 6 \mu\text{m}$ and $10 \mu\text{m}$ for $\Delta_X = 29 \mu\text{m}$. (b) The experimental measurements, indicated by a point for the mean value, with the range of values represented by the red shaded region and the 'hatched' box representing the standard deviation, compared with the range of $2\omega_L$ signals from the simulations shown in (a), represented by the green shaded region.

6.6.4 Second harmonic signal as a function of Δ_X value

Finally, this section investigates the change in $2\omega_L$ signal as a function of Δ_X . This is conducted in a similar way to previous sections, in which the results from simulations are compared with the experimental measurements. This analysis is performed for the two aperture diameters used during the experiment, $d = 6 \mu\text{m}$ and $10 \mu\text{m}$.

Results for the case of $d = 6 \mu\text{m}$

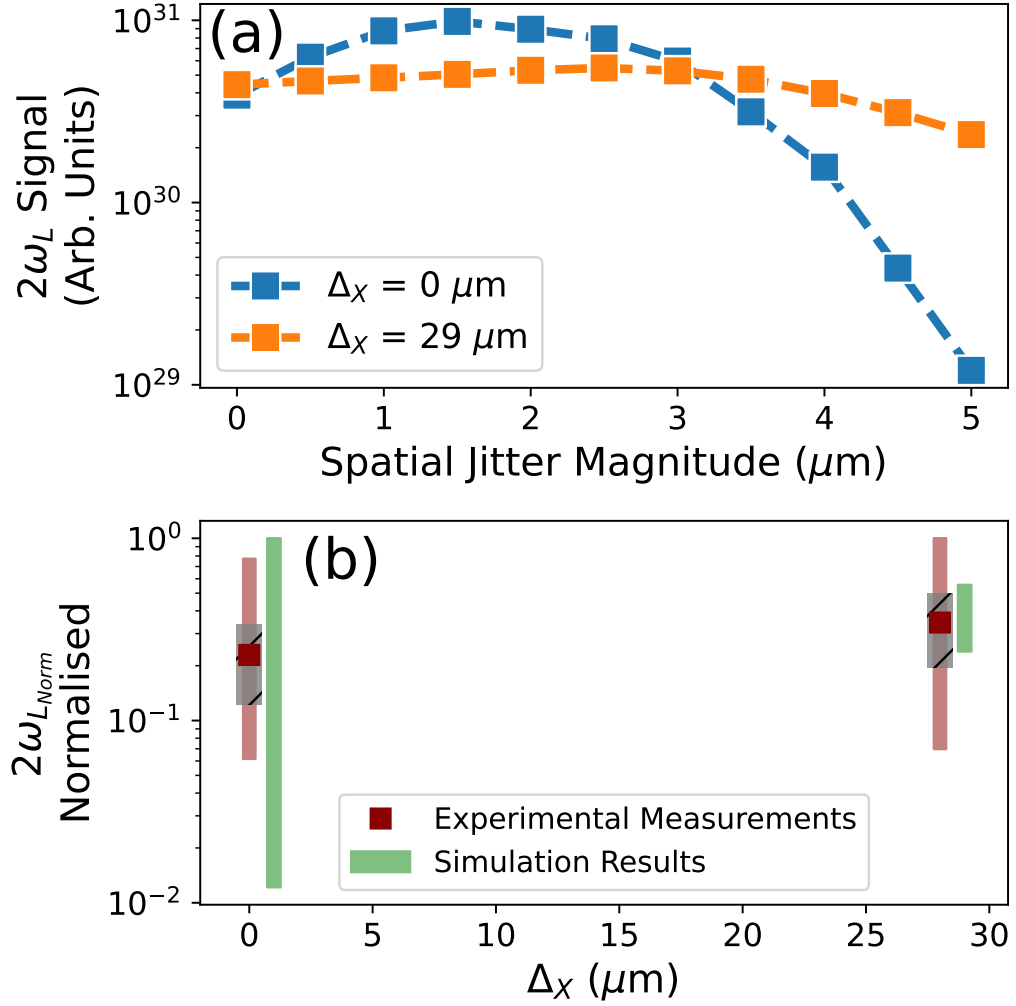


Figure 6.12: (a) Results of the two sets of simulations performed in which spatial jitter magnitude was varied for targets with $\Delta_X = 0 \mu\text{m}$ and $29 \mu\text{m}$ for $d = 6 \mu\text{m}$. (b) The experimental measurements, indicated by a point for the mean value, with the range of values represented by the red shaded region and the 'hatched' box representing the standard deviation, compared with the range of $2\omega_L$ signals from the simulations shown in (a), represented by the green shaded region.

Figure 6.12(a) shows the change in $2\omega_L$ signal as a function of spatial jitter magnitude for $\Delta_X = 0$ and $29 \mu\text{m}$, taken from two sets of simulations for a fixed value of $d = 6 \mu\text{m}$. As can be seen, both the minimum and maximum signal is produced for $\Delta_X = 0 \mu\text{m}$ for this range of spatial jitter magnitudes. Figure 6.12(b) shows the experimental measurements displayed alongside the range of the simulation results, as a function of Δ_X , displayed in the same manner as in the previous figure.

The comparison between simulation and experimental measurement demonstrate

that for the case of $\Delta_X = 0 \text{ } \mu\text{m}$, the range of experimental measurements sits within the range of $2\omega_L$ signals observed in the simulations. However, this is clearly not the case for $\Delta_X = 29 \text{ } \mu\text{m}$, with the range of experimental measurements clearly exceeding that from the simulations. In addition, the maximum signal measured during the experiment is observed for $\Delta_X = 29 \text{ } \mu\text{m}$, which is not seen in the simulation results. This again demonstrates that there is greater disagreement between the experimental measurements and simulation results for greater values of Δ_X . Again, this may be due to the non-Gaussian focal spot shape observed under these interaction conditions.

Results for the case of $d = 10 \text{ } \mu\text{m}$

Figure 6.13(a) shows the change in $2\omega_L$ signal as a function of spatial jitter magnitude for $\Delta_X = 0, 20$ and $29 \text{ } \mu\text{m}$, taken from two sets of simulations for a fixed value of $d = 10 \text{ } \mu\text{m}$. As can be seen, both the minimum and maximum signal is produced for $\Delta_X = 0 \text{ } \mu\text{m}$ for this range of spatial jitter magnitudes. Figure 6.13(b) shows the experimental measurements displayed alongside the range of the simulation results, as a function of Δ_X . The simulation results have been normalised to the experimental measurements for each Δ_X value to enable direct comparison. The range of experimental measurements are shown, displayed in the same manner as in the previous figure.

Similar to the investigation of the measurements made for $d = 6 \text{ } \mu\text{m}$, there is agreement between the range of experimental measurements and simulations results for $\Delta_X = 0 \text{ } \mu\text{m}$. The maximum signal measured during the experiment for these interaction conditions is for $\Delta_X = 20 \text{ } \mu\text{m}$, not aligning with the simulation results and the range of $2\omega_L$ signals seen for $\Delta_X = 20$ and $29 \text{ } \mu\text{m}$ exceed those seen from the simulation results. This once again suggests that the non-Gaussian focal spot spatial intensity profile may have a large impact on the generation of frequency-doubled light.

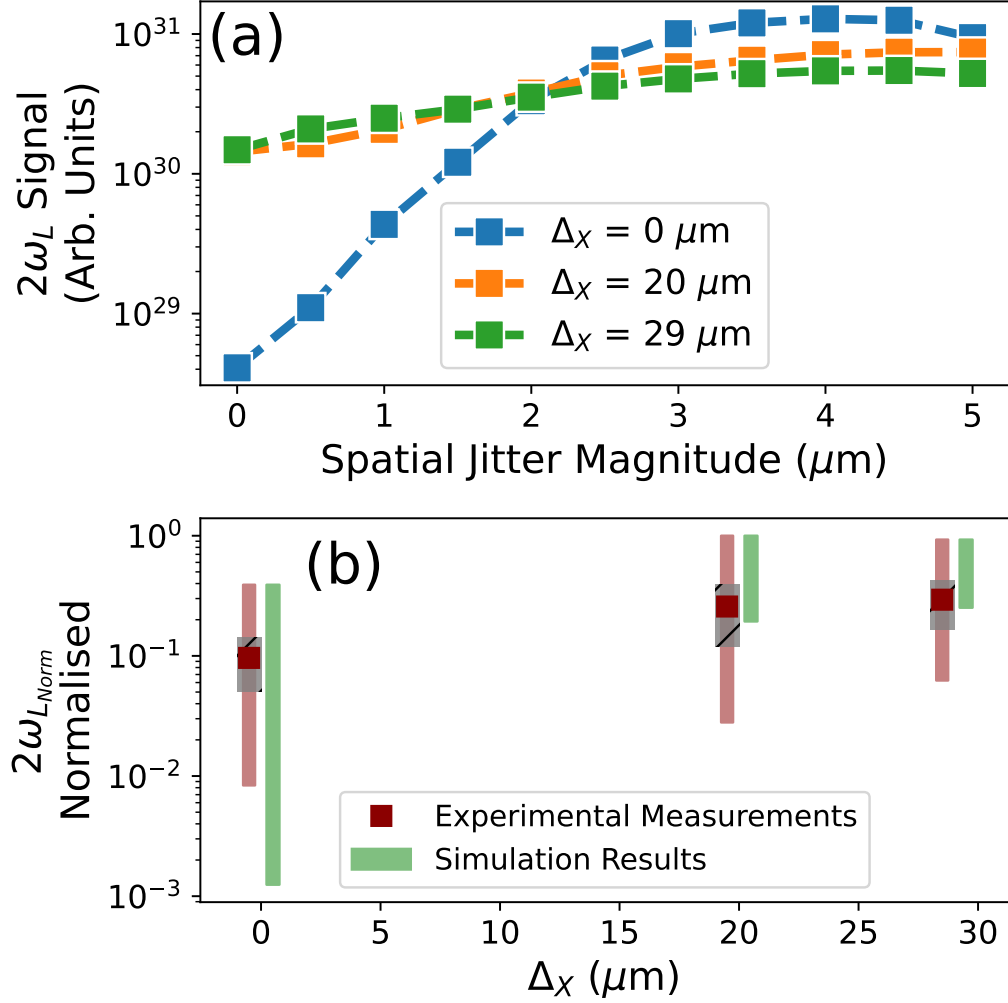


Figure 6.13: (a) Results of the two sets of simulations performed in which spatial jitter magnitude was varied for targets with $\Delta_X = 0, 20$ and $29 \mu\text{m}$ for $d = 10 \mu\text{m}$. (b) The experimental measurements, indicated by a point for the mean value, with the range of values represented by the red shaded region and the ‘hatched’ box representing the standard deviation, compared with the range of $2\omega_L$ signals from the simulations shown in (a), represented by the green shaded region.

6.6.5 Modelling the effect of a non-Gaussian focal spot

The experimental measurements discussed above have pointed to the possibility that the non-Gaussian focal spot that occurs at the point of interaction between the pulse and aperture target for $\Delta_X = 20$ and $29 \mu\text{m}$ may be responsible for the lack of agreement between experimental measurements and simulations results. Focal spot images were not captured before every shot, and as there was a degree of variability in the focal spot spatial distributions for fixed Δ_X values, it is not possible to directly account for their effect on $2\omega_L$ generation. Instead, to model the influence of a non-Gaussian focal spot spatial intensity distribution on the generation of $2\omega_L$ light, the approach was taken to

run several sets of simulations in which a non-Gaussian focal spot shape was used that is quantifiable and the energy distribution of which could be programmatically varied. These sets of simulations served to test the hypothesis that a non-Gaussian focal spot shape does in fact alter the generation of $2\omega_L$ light from an aperture target, even if the results cannot be directly compared with experimental measurements.

For these sets of simulations, the grid was defined as 3200×3200 mesh cells, corresponding to a simulation box size of $45 \mu\text{m} \times 45 \mu\text{m}$, with a particle-per-cell number of 15. The target consisted of Al^{13+} ions of a fixed thickness $l = 6 \mu\text{m}$ with a fixed aperture diameter of $d = 9 \mu\text{m}$. The ions were neutralised with an electron population with a peak density of $30n_{crit}$ and an initial temperature of 10 keV. The pulses used, entering the simulation box from the left (X) boundary, had a wavelength of $\lambda_L = 800 \text{ nm}$ and a Gaussian temporal profile with a duration of $t_L = 40 \text{ fs}$ (FWHM). The interaction was simulated for 200 fs. To create a quantifiable non-Gaussian transverse spatial profile for the pulse that could be programmatically altered, two pulses with standard Gaussian spatial profiles were propagated at the same time. Each individual pulse had a Gaussian spatial profile with a spot size of $\phi_L = 3 \mu\text{m}$, with the pulses separated by $7 \mu\text{m}$ in the Y direction, defined as the distance between the points of maximum intensity of each pulse. Figure 6.14(a) shows an example of the double-peaked spatial intensity profile this generates, somewhat similar to the spatial profile seen in figure 6.3(b) if a lineout is taken through $X = 10 \mu\text{m}$. This profile was then varied by altering the percentage of a total fixed quantity of energy present in each of the two pulses. This was expressed as a ratio of the energy present in pulse B to the total energy in both pulses, $\alpha_L = E_B/(E_A + E_B)$, with the spatial intensity profile for the case of $\alpha_L = 0.7$ shown in figure 6.14(b). The intensity of each pulse, when the energy was distributed equally, was $I_{LA} = I_{LB} = 0.5 \times 10^{21} \text{ W/cm}^2$.

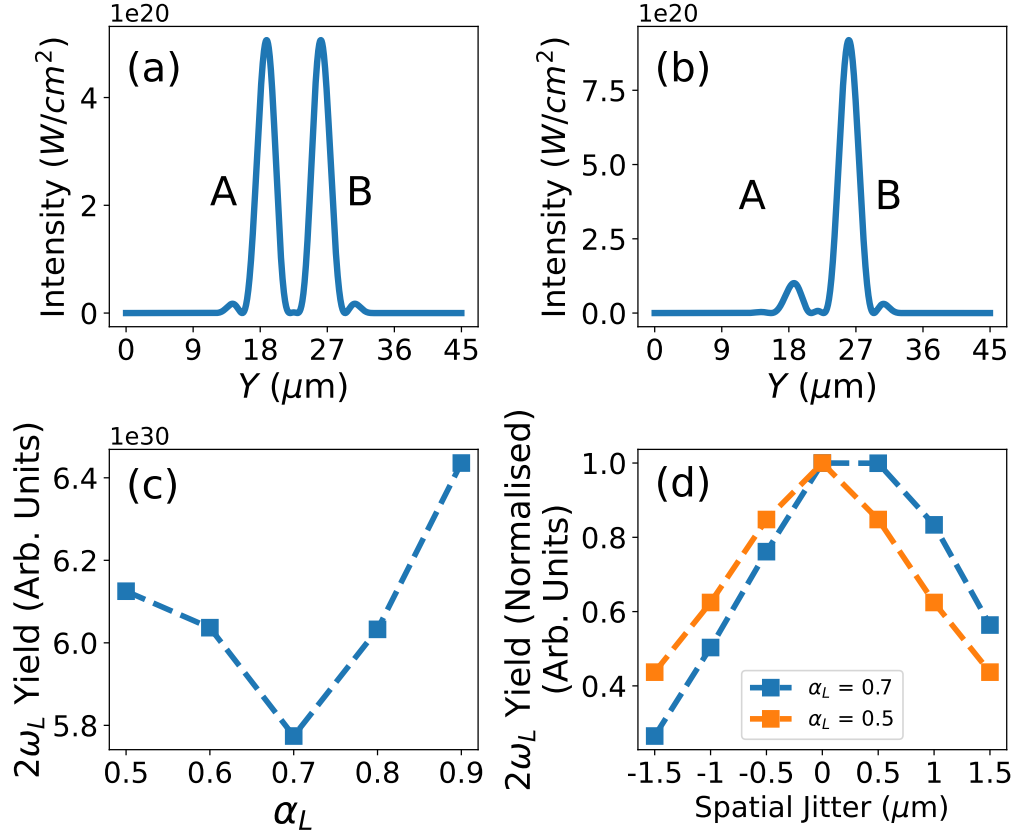


Figure 6.14: Figures (a) and (b) show line-outs taken through the focal spots used in EPOCH simulations to model the effects of a non-Gaussian focal spot on the operation of the aperture mechanism. The first shows an equal distribution of energy between the two co-propagating pulses, the second the case of 70% to 30%. (c) The effect of changing the energy distribution between the two pulses on the $2\omega_L$ yield produced during an interaction. (d) The difference between the effects of $\alpha_L = 0.7$ and $\alpha_L = 0.5$ on an interaction when the spatial jitter is varied.

The first set of simulations that were performed involved varying the distribution of energy between the pulses, α_L , from 0.5 to 0.9 and examining the variation in $2\omega_L$ yield. The results are shown in figure 6.14(c). As can be seen, changing this energy distribution results in a change in the $2\omega_L$ yield by $\pm \sim 5\%$. Figure 6.14(d) shows the results of the second set of simulations performed that involved varying the spatial jitter for two separate values of α_L , 0.5 and 0.7. The spatial jitter in this case is defined as the simultaneous movement from the center of the simulation box of both pulses that comprise the non-Gaussian spatial profile. In the case of the symmetric pulse profile, $\alpha_L = 0.5$, the $2\omega_L$ yield is also symmetric with positive and negative spatial jitter value. For the case of the asymmetric spatial profile, $\alpha_L = 0.7$, the $2\omega_L$ yield is also asymmetric with spatial jitter value. As the portion of the pulse containing the greater share of the energy interacts with one of the aperture edges (positive spatial

jitter) the $2\omega_L$ yield is greater than in the inverse case

The results of these simulations demonstrate that for a pulse with a non-Gaussian spatial profile, even for a simplified test case such as this, the distribution of energy within the pulse effects the $2\omega_L$ yield. In addition, when a non-Gaussian pulse with an asymmetric energy distribution interacts with an aperture, it is not only the spatial jitter magnitude that must be considered, as has been the case up until this point, but whether the jitter is positive or negative. To fully characterise the effects of the spatial intensity distribution on the $2\omega_L$ light generation, the exact distributions would be required. In addition, the use of 3D simulations would likely be required due to the lack of rotational symmetry displayed from the intensity distributions captured during the experiment, a factor that could not be accounted for through the use of 2D simulations.

6.7 Conclusion

This chapter reports on the results of an experimental campaign conducted using the Gemini laser system to investigate the generation of frequency doubled light from pre-formed aperture targets. This built directly off of the work presented in the previous two chapters of this thesis in which the generation of this form of light was investigated via numerical simulations. Aperture targets with two diameters were used throughout the experiment, $d = 6$ and $10\ \mu\text{m}$. The quantity of ω_L light transmitted through and $2\omega_L$ light generated from these targets was measured as a function of pulse energy and focal spot size. The effect of the focal spot spatial intensity distribution was explored through moving the target toward the focusing OAP, with three positions chosen, $\Delta_X = 0, 20$ and $29\ \mu\text{m}$. Accompanying the resultant increase in focal spot diameter was a notable change in the focal spot spatial intensity profile, away from the more Gaussian shape observed at $\Delta_X = 0\ \mu\text{m}$.

The experimental measurements were understood through comparison with the results of combined FLASH and EPOCH simulations performed under the same interaction conditions. From this, it was found that when the focal spot of the pulse was of sufficiently quality, i.e. Gaussian, and with a FWHM that did not vary greatly in size, the results largely agreed with the model put forward in the previous two chapters. Observed under these interaction conditions were a range of $2\omega_L$ signals which

could largely be explained by considering the role of the on-shot variation in spatial jitter. As the focal spot size was varied through altering the target defocus position, the measurements made could not be explained through spatial jitter effects alone. A possible reason for this was the change in focal spot spatial intensity distribution to one of a non-Gaussian nature that accompanied the increase in target defocus position. As directly accounting for the non-Gaussian focal spot spatial profile in the simulations was not possible, a simplified model of a non-Gaussian focal spot was investigated, showing that it does in fact impact the generation of $2\omega_L$ light.

With respect to future experimental investigations of aperture targets for the generation of $2\omega_L$ light, this study demonstrates that the generation of this light is likely highly sensitive to both the effects of spatial jitter magnitude and the quality of the focal spot spatial intensity profile. Mitigating the effects of the spatial jitter magnitude could be achieved through the use of a larger focal spot size. However, this chapter reveals that if this is accomplished solely by moving the target out of the focus of the laser, the potential benefits are outweighed by the focal spot spatial intensity distribution deviating from the expected Gaussian shape and the ensuing effect this has on the interaction. Should a change in focal spot size be achieved through the same method of altering the target position, steps should be taken to ensure that the focal spot shape can either be maintained or accounted for when analysing any measurements made. Alternatively, the use of a focusing optic with a greater F-number would lead to a larger focal spot size at best focus that possessed the desired Gaussian spatial intensity profile.

Chapter 7

Investigation of the spatio-temporal variations in the polarisation of light transmitted through relativistically induced plasma-apertures

7.1 Introduction

This chapter reports on the results of an experiment that builds on the work presented in M. Duff *et al.*[56] in which short duration (tens of fs) high-intensity ($\sim 10^{20}$ W/cm²) drive pulses irradiated ultra-thin targets, resulting in the formation of self-generated relativistic plasma apertures in the targets (as discussed in in chapter 4), allowing for the transmission of a portion of the drive pulse through the target. Similar to the aperture mechanism investigated in the previous three chapters of this thesis, the acceleration of electron bunches from the edges of the aperture by the pulse results in the generation of light from within the target. When measuring the polarisation state of the light at the rear of the target, a shift in the angle of linear polarisation (AOLP) relative to the drive pulse was observed. This shift was attributed to the super-position of light transmitted through the target and light generated from within the target. The magnitude of the shift in AOLP was found to be dependent on the transmitted light

percentage, attributable to the relative intensities of the two forms of light present.

The experiment reported on in this chapter involved high-intensity ($\sim 10^{19}$ W/cm²), short duration (~ 1 ps) pulses generated by the Vulcan laser system (discussed in section 3.3.2) irradiating ultra-thin targets (tens of nm). A dual-pulse configuration was used to investigate the changes to the polarisation state of the light at the rear of the target in an attempt to determine if similar changes to the measured AOLP can be observed when picosecond duration pulses are used, compared to those with tens of femtoseconds durations used in reference [56].

7.2 Experimental method

7.2.1 Beam-line

The experimental measurements presented in this chapter were obtained through the use of the Vulcan laser system delivering laser pulses to one of its interaction areas, known as Target-Area-West (TAW) (section 3.3.2). Of the eight beam-lines available in TAW, the two short-pulse beams were utilised, referred to as beam 7 (B7) and beam 8 (B8). B8 pulses were employed as a ‘pre-ioniser’ used to induce pre-ionisation of the target and to control the extent of the resulting plasma expansion, prior to the arrival of the main-beam B7 pulses. B7 delivered pulses with minimum durations of $t_{L7} \sim 1$ ps (FWHM) and energies up to $E_{L7} \sim 100$ J. B8 delivered pulses of the same duration while being apodised to a diameter of 25 mm to reduce its energy content to a value of up to $E_{L8} \sim 1.5$ J and to create a larger focal spot relative to that of B7 pulses. This was desirable as it would reduce the possibility of the two focal spots not overlapping spatially during the interaction. Both pulses have a wavelength of $\lambda_L = 1053$ nm. To ensure that pulses from both beam-lines followed the same path prior to interaction with the target, B8 pulses were injected through a hole in one of the mirrors used to guide B7 pulses. The total path lengths of the two beam-lines in use could be configured such that B8 pulses could interact with the target either prior to, or after, the arrival of the higher intensity B7 pulses.

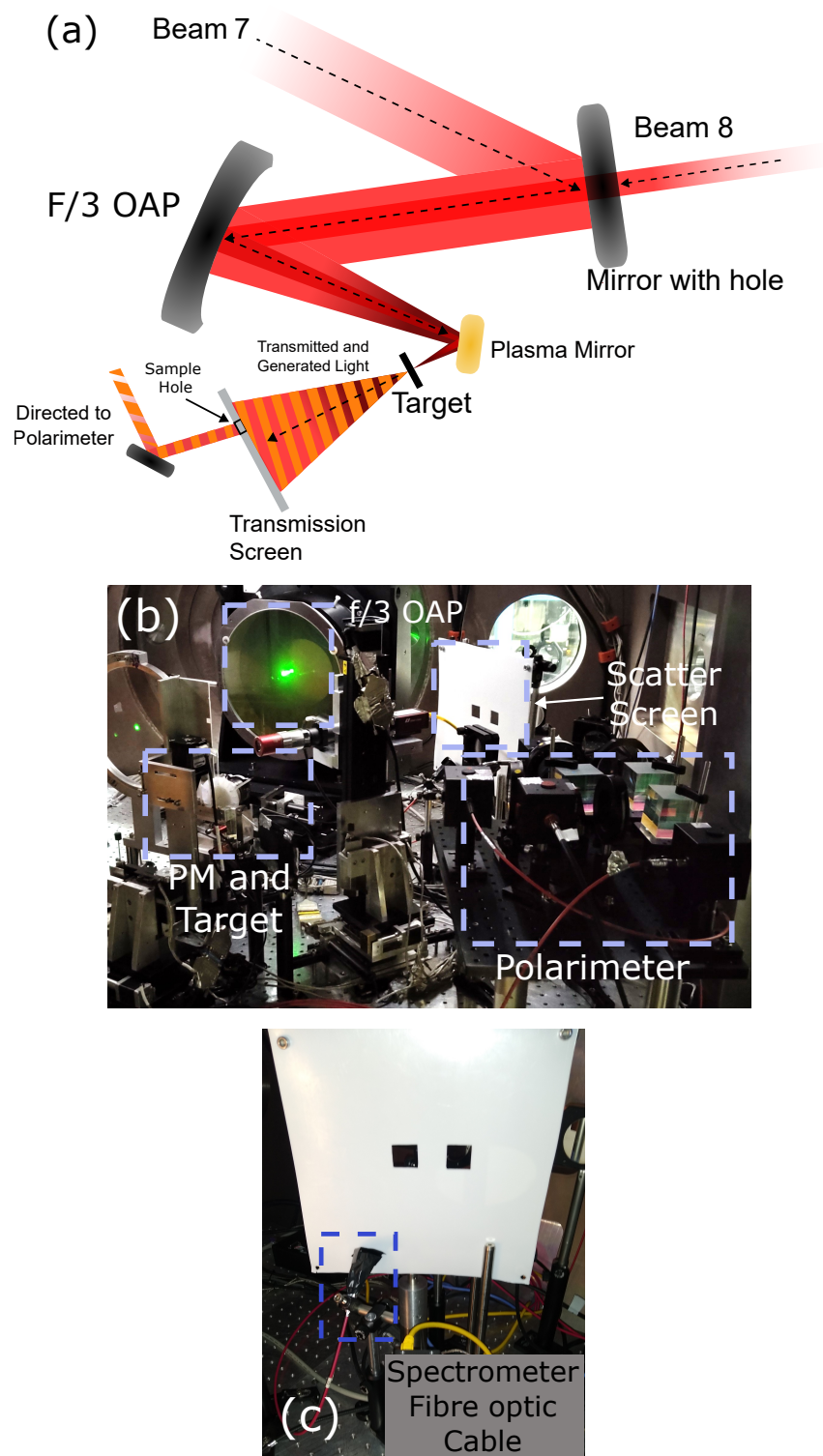


Figure 7.1: (a) Schematic of the beam-line used during the experiment. The lower intensity beam 8 is shown being injected into the main beam-line through a hole in one of the mirrors prior to the $f/3$ focusing optic. Both beams interact with PM prior to interaction with the target. The generated and transmitted light is then incident on a scatter screen for the purposes of imaging. The ‘sample hole’ in the transmission screen that allows for a portion of the light to enter the polarimeter is shown. (b) An image taken during the experiment of the interior of the interaction chamber showing the focusing OAP, PM, target, scatter screen and polarimeter diagnostic (discussed below). (c) An image taken during the experiment showing the transmission screen and the position of the fibre optic cable connected to the optical spectrometer.

An F/3 off-axis parabola (OAP) was used to bring both beams to focus onto plastic (CH) targets of thickness ranging from $l = 10$ nm to $l = 125$ nm. B7 pulses achieved maximum pulse intensities of $I_{L7} \sim 2 \times 10^{19}$ W/cm² and those from B8 achieved maximum pulse intensities of $I_{L8} \sim 1 \times 10^{14}$ W/cm². An anti-reflection (AR) coated plasma mirror (PM) was employed to increase the pulse contrast by up to two orders of magnitude to create more favourable interaction conditions. The PM was placed such that it reflected both beams onto the target. As only B7 generated pulses with sufficient intensity to activate the PM, pulses from B8 did not experience an increase in temporal intensity contrast but were attenuated by the AR coating of the PM by roughly two orders of magnitude. Note that the PM is reflective in its non-plasma mode, allowing B8 to reach the targets.

A schematic of the beam-line displaying the elements mentioned above is shown in figure 7.1(a) in addition to several photographs taken during the experiment of the beam-line and diagnostic components in (b) and (c), respectively.

7.2.2 Diagnostics

To investigate the light generated from and transmitted through the target, three diagnostics were employed: a transmission diagnostic to image the spatial-intensity distribution of the light at the rear of the target, an optical spectrometer to measure its spectrum over a wavelength range of 500 to 1200 nm, and a polarimeter to characterise its polarisation state. For the transmission diagnostic, a diffuse scatter screen was positioned 50 cm from the rear of the target, allowing the fully expanded beam to be incident upon it. A CCD camera, filtered for ω_L frequency light, with a $1.2\times$ telephoto lens was used to image the light incident on the scatter screen. For spectral measurements, a fibre optic cable was placed near the transmission scatter screen, facing the center, connected to a spectrometer located outside the interaction chamber. The polarimeter, details of which are provided below, was placed within the interaction chamber. This step was taken as the light being measured was brought to focus before entering the polarimeter. If the polarimeter were placed outside the chamber, there was the possibility that as the light was brought to focus, it would be of sufficient intensity to ionise the air. The resulting change in the optical properties of the air could affect the path of the light such that it did not reach the polarimeter or it could alter its polarisation, thus affecting the measurements made with this diagnostic.

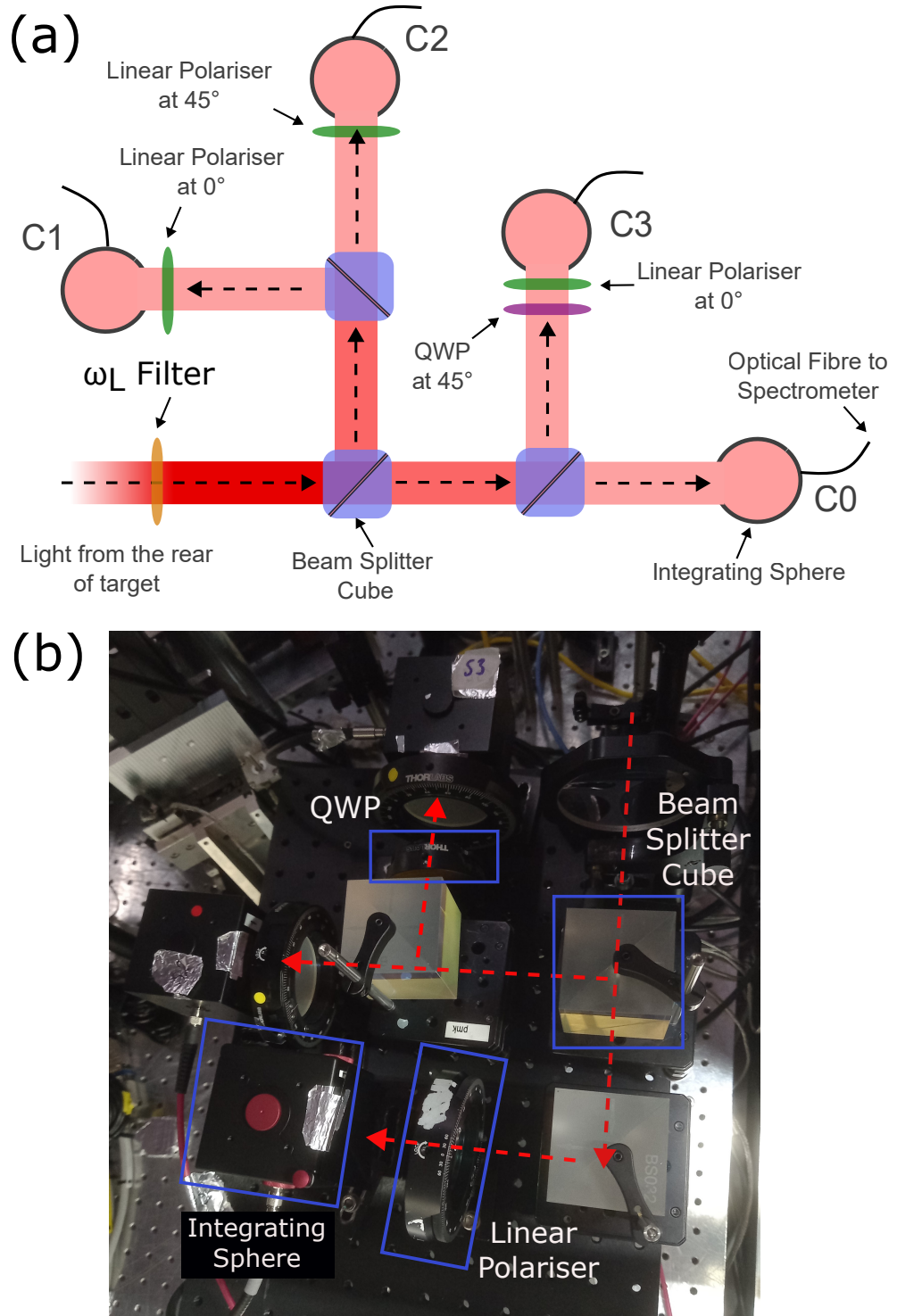


Figure 7.2: (a) Illustration of the polarimeter setup used to measure the polarisation state of light from the rear of the target. The signal from each channel, C_n corresponds to the Stokes parameter S_n in equations 7.1 to 7.4. (b) A photograph taken during the experiment of the polarimeter with the various components labelled.

A sub-sample of the light from the rear of the target, roughly 4%, was measured by the polarimeter by cutting a small sample hole into the transmission screen, behind which a mirror was placed that directed this light into the polarimeter. This allowed for the simultaneous use of the both the transmission diagnostic and polarimeter. Figure 7.1(a) illustrates the positioning of the scatter screen and the mirror used to direct light into the polarimeter. An illustration of the polarimeter diagnostic setup and a photograph taken during the experiment of the polarimeter as it was positioned within the interaction chamber are shown in figures 7.2(a) and (b), respectively. As described in section 3.4.3, to enable the measurement of the polarisation state of light, the polarimeter splits a sample of light into four channels, with measurements from each channel, C_n , then used to calculate a corresponding Stokes component, S_n , and thus a Stokes vector. From this vector the AOLP of the light being measured can be calculated.

As seen, the incoming light, filtered for ω_L , is divided using beam-splitter cubes (BS), that evenly split the light into two separate paths without altering its polarisation. The light in each channel is then collected using an integrating sphere (IS), a component that consists of a hollow sphere, the interior of which is coated in a material that uniformly scatters any light within it. The spatial information of the light that enters the IS is lost, but the quantity of the light is measurable. In each path, a fibre optic cable is used to direct the light from the IS to an optical spectrometer. From the spectrometer trace measured on each channel, a single value, C_n , is extracted and used with the following equations to calculate the AOLP.

$$S_0 = C_0 \quad (7.1)$$

$$S_1 = \frac{2C_1}{S_0} - 1 \quad (7.2)$$

$$S_2 = \frac{2C_2}{S_0} - 1 \quad (7.3)$$

$$S_3 = \frac{2C_3}{S_0} - 1 \quad (7.4)$$

$$\text{AOLP} = \frac{1}{2} \arctan\left(\frac{S_2}{S_1}\right) \quad (7.5)$$

7.3 Experimental results

Analysis of transmitted light images

To obtain a measure of the total photon count of the transmitted and generated ω_L light incident on the scatter screen, the method employed was to sum the signal from an image captured by the CCD camera. As with the experimental measurements reported in the previous chapter (6.3), scattered light from other objects within the chamber is present in the images. This necessitated the cropping of the images to remove the unwanted spurious signal. An example of a transmission screen image before and after the cropping process is shown below in figures 7.3(a) and (b), respectively.

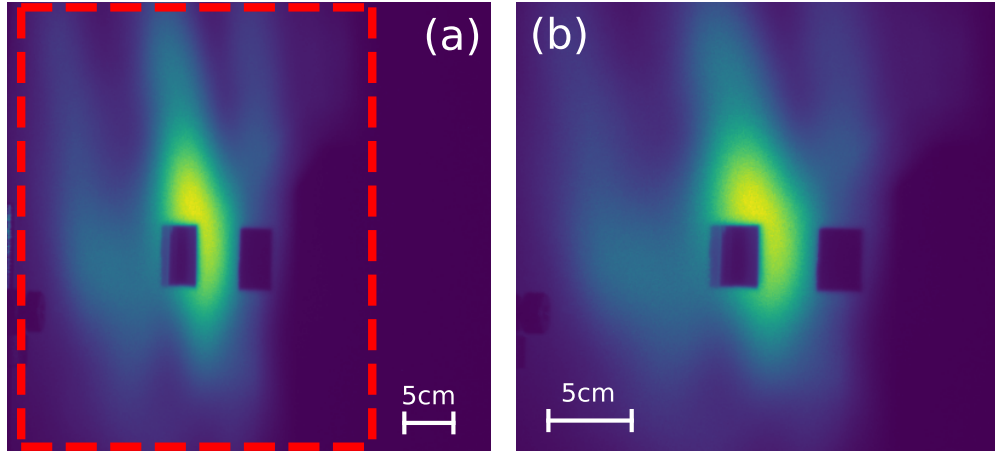


Figure 7.3: Example of an imaged capture by the CCD camera of the transmission screen, filtered for ω_L light (a) before and (b) after the cropping process.

Minor movements in either the scatter screen position or CCD camera being used to capture the images over the course of the experiment lead to the preserved area of the image being adjusted when necessary. Note that there are two holes present in the transmission screen. These were required in order to move the mirror used to reflect light into the polarimeter to a different position behind the transmission screen at various points during the experiment to allow for the placing of other diagnostics.

7.3.1 Drive pulse transmission

As discussed in the introduction to this chapter, M. Duff *et al.* [56] demonstrated that the measured AOLP of the light at the rear of the target was found to vary with the transmitted light percentage. As such, the first stage of this experiment involved determining a target thickness that would allow for a suitable range of transmitted light percentage values to be observed. This was done by conducting a series of shots in which B7 pulses were interacted with targets with thicknesses ranging from $l = 10$ to $l = 125$ nm, with this range selected based on previously published transmission measurements under similar interaction conditions [75][110][111]. The aim was to find a target thickness for which the transmission percentage was in the range of 30-40% which, when combined with an alteration of the B7 pulse energy or a change in the delay between the B8 and B7 pulses, would likely provide transmission percentages in the range of ~ 10 to 60%, potentially allowing for the measurement of a range of AOLP of the light at the rear of the target.

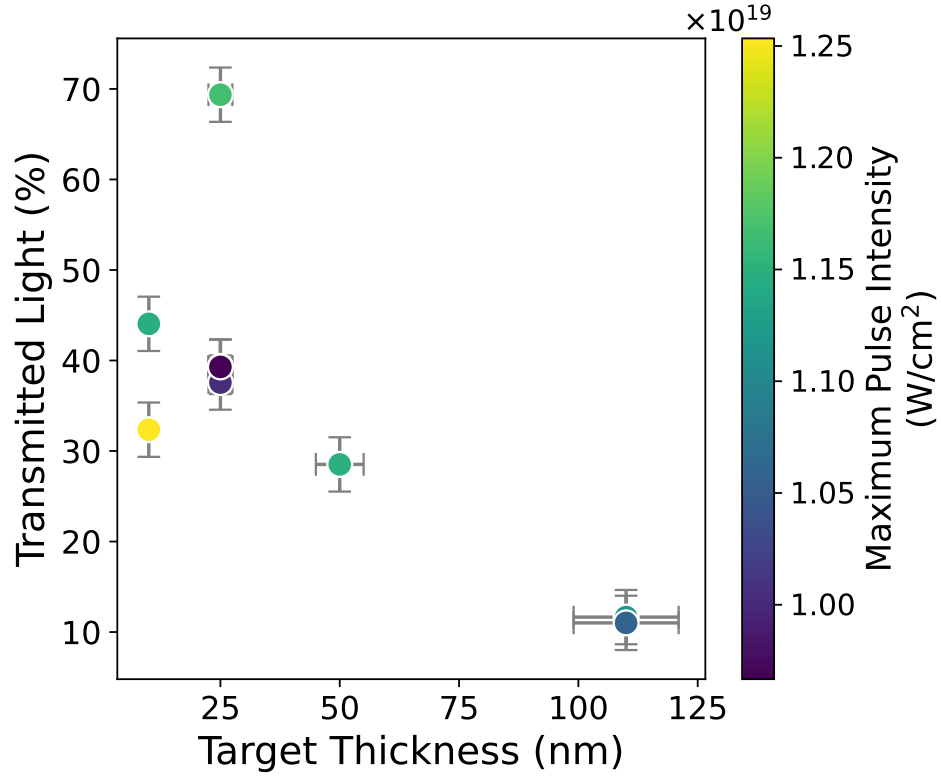


Figure 7.4: The percentage of the B7 pulse transmitted through the target as it varies with target thickness. The colour bar represents the maximum pulse intensity for each shot.

To correctly characterise the percentage of transmitted light through the target, a no-target shot was conducted at 60% of the maximum available pulse energy. This value was chosen based on previous work in which the transmitted light percentage for a similar target thickness range was investigated [75]. The purpose of this no-target shot was to emulate a shot during which the target undergoes transparency as early as possible during the interaction, and is therefore the maximum possible quantity of light that could be transmitted through the target, allowing for a comparable measurement to be used to calculate the percentage transmission and for the correct neutral density filtering to be used.

Figure 7.4 shows the transmitted light percentages measured for the range of target thicknesses. The maximum B7 pulse energy was requested for all shots, giving values in the range $E_{L7} = (99 \pm 9)$ J. The uncertainties in the target thickness are quoted from the manufacturer and those for the transmitted light percentage arise from calibrating the light level on the transmission screen. The trend observed is largely in agreement with previous published results under similar interaction conditions, aside from the near 70% transmitted light percentage measured for an $l = 25$ nm target. To aid in determining whether this was an anomalous measurement, the calculated maximum pulse intensity for each shot are included in the colour bar of this plot. Examining the intensity of the outlying measurement, it is observed to be $\sim 20\%$ higher than the other shots taken with a 25 nm target. This increased intensity would cause the target to undergo transparency earlier in the interaction, explaining the higher transmitted light percentage for this shot.

Based on the data displayed in figure 7.4, a target thickness of 25 nm was selected as it would likely allow for the exploration of a relatively large parameter space within which a range of polarisation states of the light at the rear of the target could be detected.

7.3.2 Transmitted light with varying pre-ioniser delay

With a target thickness chosen, the next step was to examine how the arrival of B8 pulses prior to the main pulse influenced the transmitted light percentage. Acting as a pre-ioniser, the B8 pulses were intended to control the degree to which the main pulse was transmitted through the target by varying the delay between the two pulses. Increasing the delay would allow more time for plasma expansion, which would be

expected to lead to transparency to B7 pulses occurring earlier in the interaction. This would result in a higher percentage of the main pulse being transmitted through the target.

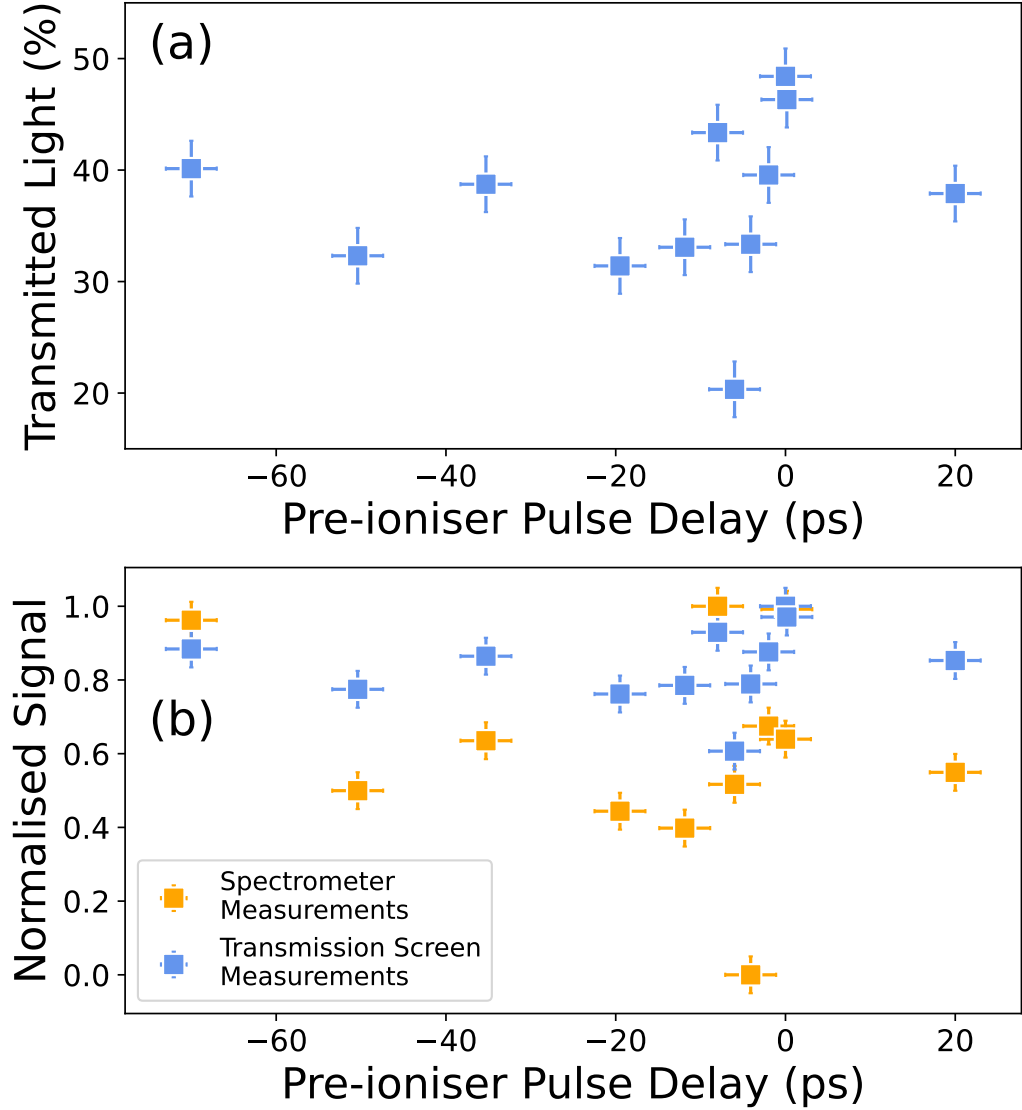


Figure 7.5: (a) Change in the percentage of light from B7 pulses that was transmitted through the target as a function of pulse delay between the pre-heater and drive pulse. (b) The transmitted light percentage and the spectrometer measurements as a function of pre-ioniser delay, normalised to allow for direct comparison. All measurements were made for targets of $l = 25$ nm, $E_{L7} = (103 \pm 8)$ J and $E_{L8} = (1.2 \pm 0.1)$ J.

This was tested through a series of shots in which the delay between the two pulses was varied from -70 to $+20$ ps, with a negative value indicating that a B8 pulse arrived before that from B7. Positive delay values were included to test the assumption that this would have no impact on the interaction as the contribution to the total quantity

of measured light from the B8 pulses that would be transmitted in this case would be negligible. The effect of this on the transmitted light percentage is shown in figure 7.5(a).

The transmission percentage observed in the results does not follow the expected trend of greater transmitted light percentage with increasing pulse delay. Instead, what is observed is a somewhat random variation with pulse delay. Unlike the clearer relationship seen between transmission percentage and target thickness, shown in figure 7.4, this is not due to a single or a few outlier data points. Although it is highly unlikely for a CCD camera diagnostic to produce 12 spurious signals, a simple verification was performed to rule this out. This involved examining the change in the summed ω_L transmission spectrometer data, normalised to pulse energy, as a function of pulse delay. The spectrometer data was summed by selecting a range of wavelength values, ± 30 nm, around the ω_L frequency and calculating the total signal within this range. This is shown in figure 7.5(b), in addition to the transmitted light percentage, both normalised to enable direct comparison. The trends seen in this figure are similar to those seen in figure 7.5, suggesting that measured transmitted light percentages are not spurious. Note that the spectrometer is not spatially integrating all of the light in the same fashion as the measurements made by the CCD camera, so identical trends would not be expected.

A possible explanation for these measurements is the influence of phenomenon such as self-focusing [102] and laser filamentation within the plasma [112] on transparency time. Frazer *et al.* [103] reported increases in pulse intensities due to self-focusing using pulses with a similar duration both experimentally and in simulations. Reported in Gray *et al.* [113] is an optimal scale length at the target front that results in increased energy absorption due to self-focusing and channelling and that for non-optimal scale lengths, laser filamentation resulted in reduced energy absorption. It is therefore possible that the conditions created by the B8 pulse at the target front resulted in self-focusing, leading to transparency occurring earlier in the interaction than would be expected, or that energy absorption was reduced such that transparency occurred later than expected. These factors may have created a more complex relationship between the pre-ioniser pulse delay and transmitted light percentage than was predicted.

Testing for activation of the plasma mirror by B8 pulses

A final check performed was that of the possible activation of the PM by B8 pulses. If this did occur, the wavefront quality of a B7 pulse would be affected and the percentage of light reflected by the PM could be significantly altered. A no-target shot with only a B7 pulse was compared with an additional no-target shot involving both beams and a pulse delay of 50 ps. This delay was chosen as there would likely be a degree of plasma expansion significant enough to alter the wavefront of the B7 pulse if activation of the PM by the B8 pulse did occur. Figure 7.6 shows the transmission screen images captured for both of these shots.

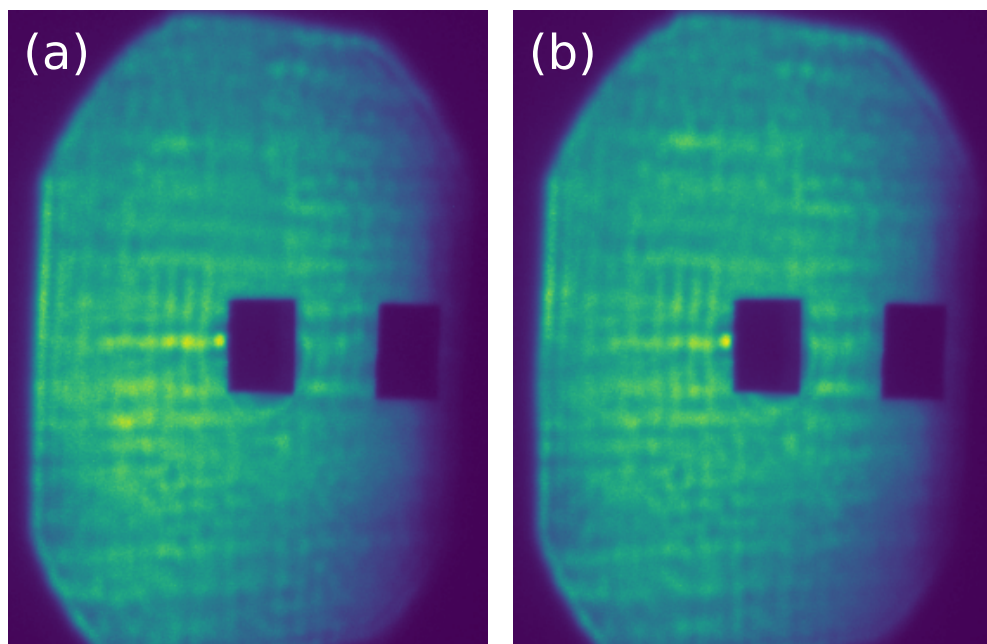


Figure 7.6: Transmission ω_L screen image taken for a shot in which (a) only a B7 pulse was used and (b) both B7 and B8 pulses were used with a delay between them of 50 ps.

As can be seen, the images are largely similar, showing that the wavefront of the B7 pulse was maintained and that the PM was not activated by the B8 pulse. In addition, a shot was taken for which only a B8 pulse was used and the PM visually inspected, showing no signs of activation. This gave further evidence that B8 pulses alone did not activate the PM and could not be the cause of the unexpected variations in transmitted light percentage with varying pulse delay. It is possible that self-focusing or filamentation within the plasma at the target front could alter the time at which transparency occurs. This may mask the expected difference in the transmitted light

percentage with pre-ioniser delay.

7.3.3 AOLP calculation method

Before experimental measurements of the angle of linear polarisation (AOLP) of the light at the rear of the target are discussed, the method by which the spectra from the polarimeter diagnostic were used to calculate the AOLP will be covered. Two possible methods existed, the first involved taking the maximum signal around ω_L from each trace (termed method A), while the alternative (termed method B) involves summing the spectral signal within a defined range around the wavelength at which the peak signal was observed. Figure 7.7 shows an example of the spectra measured by each spectrometer in the polarimeter for a target shot from the experiment. When the two possible methods are used to calculate the AOLP from this data, method A gives a value of 40° , whereas method B gives a value of 12° .

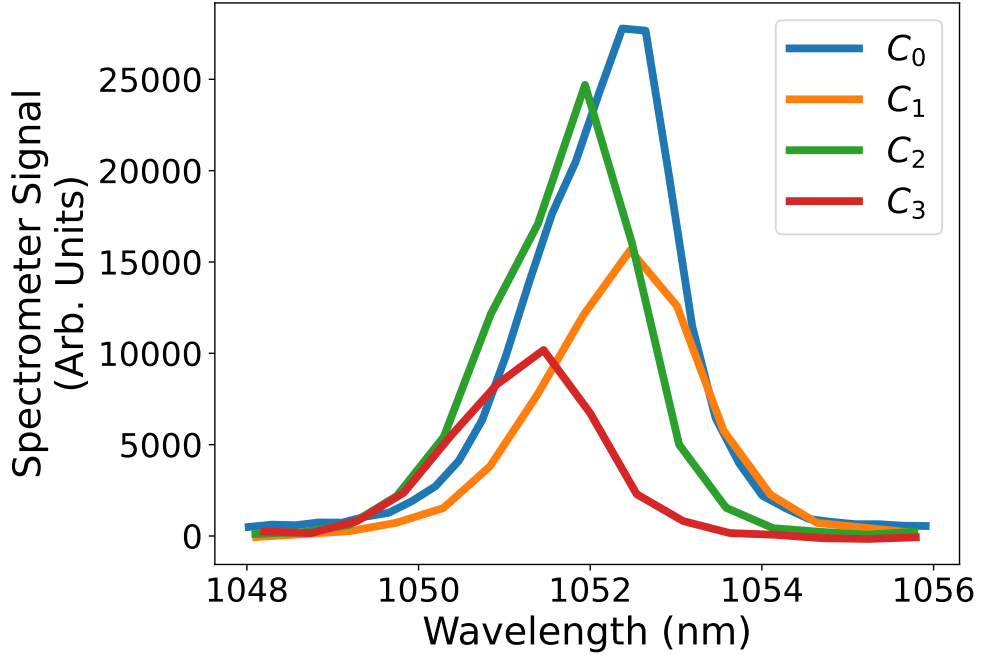


Figure 7.7: An example of spectrometer traces measured by the polarimeter for a shot involving a target with $l = 25$ nm and a pulse delay of 36 ps.

To attempt to determine which of these methods would provide the most accurate calculated AOLP, the traces from the C_1 channel of the polarimeter for measurements made for a series of shots taken using an $l = 25$ nm target during which the delay between pulses from B7 and B8 was varied were examined. These are shown in figure 7.8. Note that due to the fact that the exact wavelength at which the maximum signal is

measured, and the magnitude of this maximum value varies between shots, the step was taken to normalise their magnitudes and align all traces around their maximum values, hence the x -axis label of λ_{max} . The colour-bar values represent the transmitted light percentage measured concurrently with the spectrometer trace, included in this plot to determine if a correlation between this metric and the features of the spectra trace. Examination of the figure shows no obvious correlations between the trace shapes and transmitted light percentage, although as the trace shape is not an easily quantifiable metric, this is a qualitative judgement that is open to error.

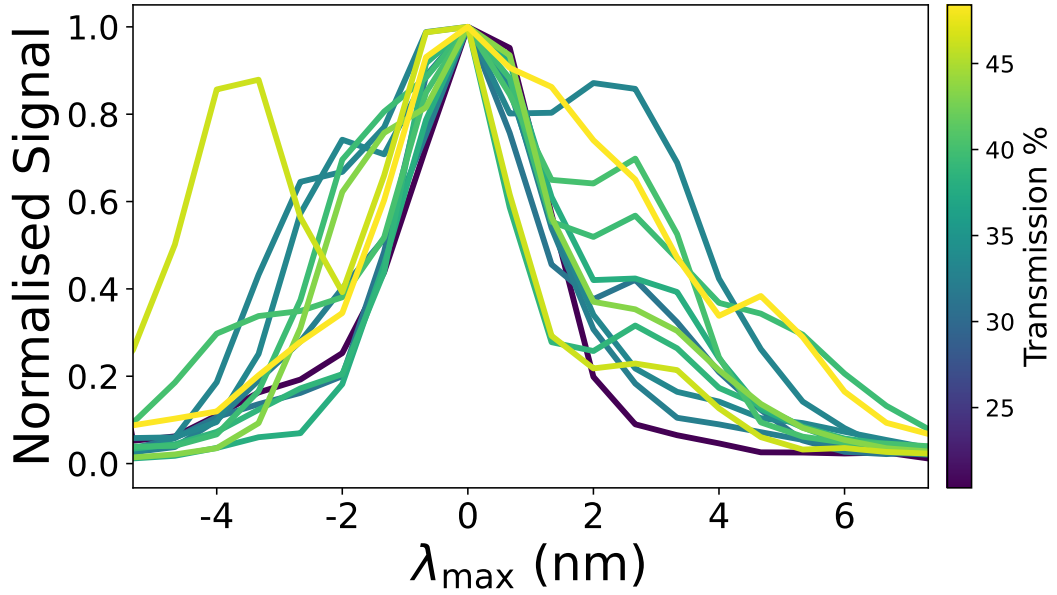


Figure 7.8: Spectrometer traces measured on the C_1 polarimeter channel for a series of shots taken during which the pulse delay between beam 7 and beam 8 was varied. The colour-bar illustrates the transmitted light percentage measured concurrently with the trace.

This probable lack of correlation seen in figure 7.8 suggest that the uneven bandwidths and asymmetries are not correlated with the AOLP and are likely diagnostic artefacts. As method B is influenced by these features of the data, the AOLP calculated using this method may not be reliable. For this reason, method A was used in the analysis of the measurements made during the remainder of the experiment, covered in subsequent sections.

Before continuing, it is worth noting a feature in figure 7.7; the wavelength at which the maximum signal is observed differs by up to 2 nm between the individual spectra measured for the four polarimeter channels. To determine whether this wavelength shift is caused by a physical process related to the interaction or another factor, such

as spectrometer calibration, a polarimeter measurement was examined for a no-target shot. Figure 7.9 shows the spectra for such a shot in which differences in the wavelength of the maximum signal are again observed but with slightly different maximum values. For example, in figure 7.7 the maximum signal for the C_0 channel is measured at a wavelength of ≈ 1053 nm, but in figure 7.9, the maximum signal for the C_0 channel is measured at a wavelength of ≈ 1052 nm. A potential cause of this could be shot-to-shot variations in the angle of the light as it enters the spectrometer or its wavefront. Both would have the effect of slightly changing the angle at which the light is reflected off of the diffraction gratings within the spectrometer [114]. This would then lead to a small shift in the apparent wavelength that has been detected.

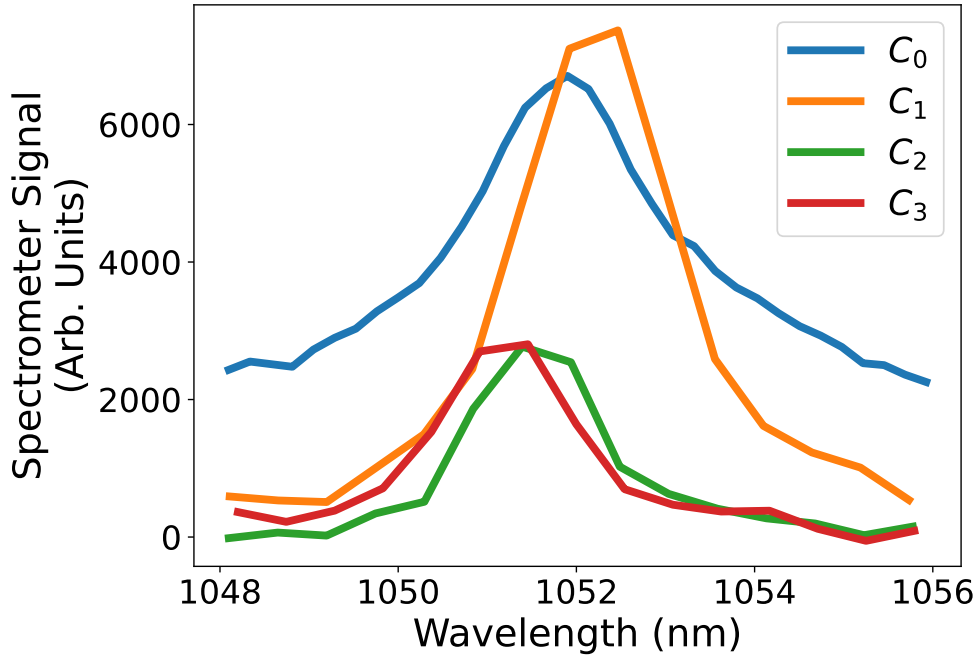


Figure 7.9: An example of spectrometer traces measured by the polarimeter for a no-target shot from the experiment, corresponding to a AOLP of 0° .

Of note in figure 7.9 is the fact that the spectrometer trace measured for the C_0 channel is not seen to drop to 0 as is the case for the other channel traces. This is due to the width of the trace for C_0 in this case being much greater than the other traces for this particular shot.

Calibration of polarimeter measurements

Present in many of the polarimeter measurements is an unexpected feature; the channel signal with the greatest value is not C_0 , with an example seen in the no-target

shot measurements shown in figure 7.9. Since the C_0 channel contains no polarisers or wave-plates that could attenuate the light, all other channel signals should theoretically be less than or equal to that of C_0 . This is likely attributable to factors such as differences in the transmission efficiencies of the fibre optic cables connecting the integrating spheres to the spectrometers, or slight variations in spectrometer sensitivities. As these factors would very likely lead to incorrect AOLP values, it was necessary to calibrate the signals measured on each channel. This was performed using a no-target shot with a known polarisation (p-polarised), shown in figure 7.9, and comparing the measured values with their expected values. In the case of p-polarised light these would be the C_0 and C_1 signals being equal in magnitude and C_2 and C_3 being half that of C_0 . Using these values, a set of calibration factors was derived. These factors were determined such that when the no-target shot channel signals were multiplied by their respective calibration factors, they aligned with the expected values for p-polarised light. These calibration factors were then applied to subsequent target-shot measurements to correct for any potential irregularities introduced by the experimental setup.

7.3.4 Change in AOLP with transmission percentage

With the verification of the method used to calculate the AOLP complete, attention can now turn to examining the polarimeter measurements made during the experiment. Although the unexpected transmission percentages observed with varying pulse delay have not been explained, a change in the polarisation state at the rear of the target would still be anticipated across the range of transmission percentages measured; 20% to 50%.

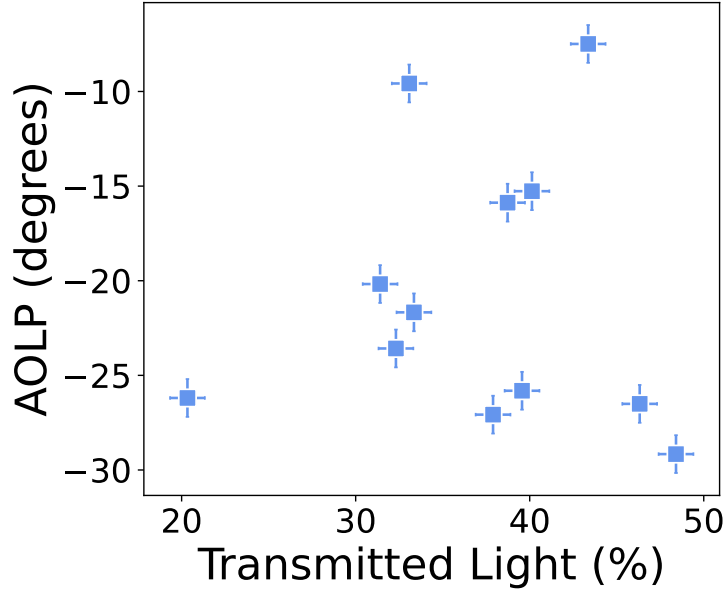


Figure 7.10: AOLP as a function of the delay between B7 and B8 pulses for targets with $l = 25$ nm. Pulse energies of $E_{L7} = (103 \pm 8)$ J and $E_{L8} = (1.2 \pm 0.1)$ J were used.

Figure 7.10 shows the relationship between the AOLP and transmitted light percentage as a result of varying the delay between B8 and B7 pulses. All measurements in this dataset were conducted using $l = 25$ nm targets, with the maximum B7 pulse energy requested, giving values of $E_{L7} = (103 \pm 8)$ J. B8 was operated at 80% of its maximum possible energy, yielding pulse energies of $E_{L8} = (1.2 \pm 0.1)$ J. Unlike in previous studies in which the change in AOLP was shown to decrease with increased light percentage, no such correlation is observed from the experimental measurements. Instead, the AOLP is seemingly random, which may be attributable to the relatively small number of measurements made. Explanation of these results is the focus of the remainder of this chapter.

7.4 Testing the operation of the polarimeter

While the method for calculating the AOLP from the polarimeter measurements had been tested, a potential source of error in the experimental measurements of the AOLP lay in the operation of the polarimeter. Although similar polarimetry diagnostics have been successfully employed in previous experimental campaigns [56], one in which the signal on each channel was measured using IS and spectrometers had not. As such, it was deemed worthwhile to determine whether the diagnostic design had any effect on

the experimental measurements. This was tested by constructing two models, one a physical replica and the other an analytical model. The channel signals for a range of polarisation states were then measured from each model and compared, with agreement between the two confirming that the diagnostic was functioning as expected.

7.4.1 Analytical model of the polarimeter diagnostic

Using the python package ‘py-pol’, an analytical model of the polarimeter was constructed in which the light being input into the polarimeter was represented by a Stokes vector and the optical components such as the polarisers and wave-plates that comprised the polarimeter were represented by Mueller matrices. The effect on light of a given polarisation as it passes through each of the polarimeter channels is then calculated using linear algebra operations. An example of such an operation representing the passing of p-polarised light through the C_2 channel is shown in figure 7.11. The input light (four-component Stokes vector on the LHS) with a normalised intensity passes through a linear polariser set to +45 degrees (Mueller matrix on the LHS), resulting in 45 degree polarised light with a reduction in intensity of 50 % (four-component Stokes vector on the RHS).

$$\begin{array}{ccc}
 \text{Linear Polariser} \\
 \text{at } 45^\circ \\
 \frac{1}{2} \begin{bmatrix} 1 & 0 & 1 & 0 \\ 0 & 0 & 0 & 0 \\ 1 & 0 & 1 & 0 \\ 0 & 0 & 0 & 0 \end{bmatrix} & \begin{bmatrix} 1 \\ 1 \\ 0 \\ 0 \end{bmatrix} & = \frac{1}{2} \begin{bmatrix} 1 \\ 0 \\ 1 \\ 0 \end{bmatrix} \\
 & \underbrace{\hspace{1cm}} & \underbrace{\hspace{1cm}} \\
 & \text{P-pol Light} & 45^\circ\text{-pol Light}
 \end{array}$$

Figure 7.11: Example of a mathematical operation used in the analytical model of the polarimeter representing the C_2 channel. Linear algebra operation between a Stokes vector representing p-polarised light propagating through a linear polariser set to 45° , represented by a Mueller matrix, with the resulting light have an AOLP of 45° and a reduction in intensity of 50%.

7.4.2 Physical replica of the polarimeter diagnostic

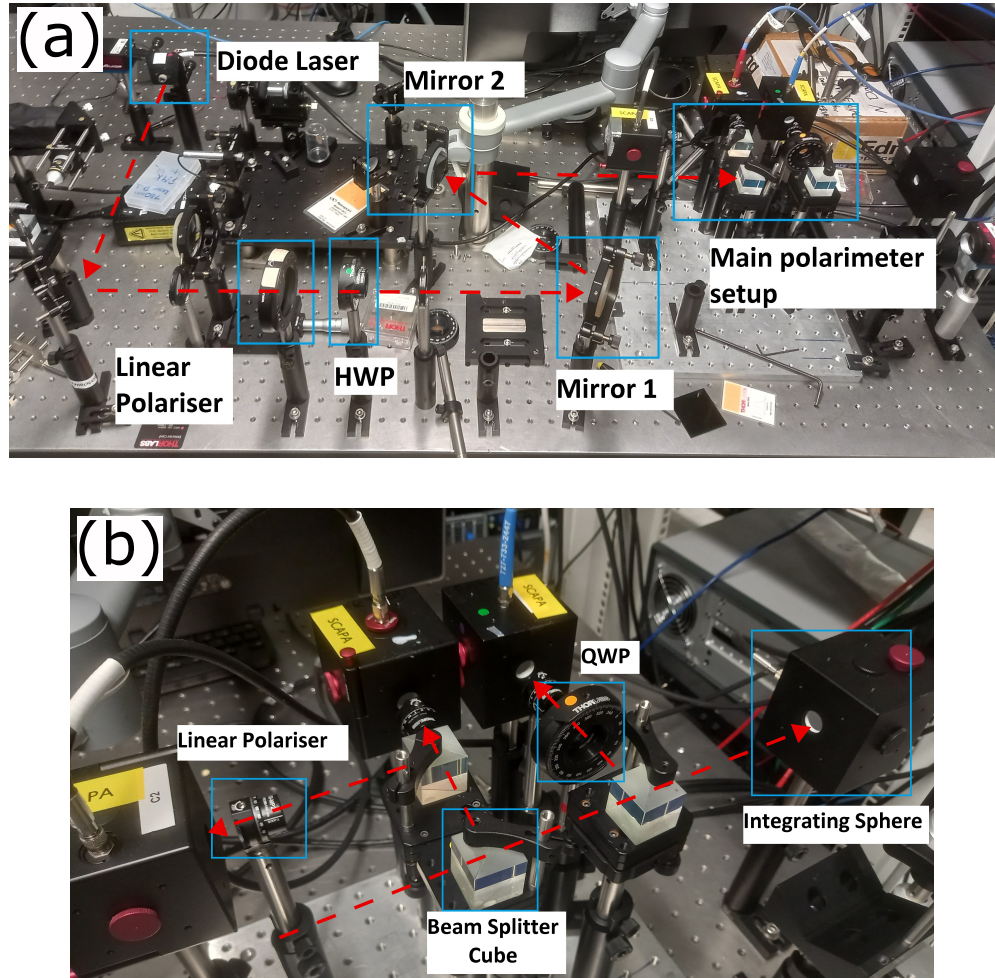


Figure 7.12: (a) A photograph of the setup used to test the response of the polarimeter when the input polarisation was altered. (b) Close-up photograph of the polarimeter setup being tested.

The physical replica of the polarimeter, constructed at the SCAPA facility at the University of Strathclyde, is shown in Figure 7.12. To generate polarised light for input into the polarimeter, a diode laser was passed through both a linear polariser and a half-wave plate (HWP), as shown in figure 7.12(a). The experimental setup for the diagnostic involved reflecting the light from the rear of the target off two mirrors before entering the polarimeter. Although these components are unlikely to affect any measurements made, they were included in the physical model for consistency, as also

shown in figure 7.12(a). Figure 7.12(b) provides a detailed view of the polarimeter components, including the beam-splitter cubes, polarisers, quarter-wave plate (QWP), and integrating spheres.

7.4.3 Comparison between polarimeter models

With the two models of the polarimeter constructed, the response of the individual channels of each to light with varying polarisation states could be measured and compared, beginning with altering the AOLP. In the analytical model, this variation is straightforward to implement, while in the physical model, it was achieved by rotating the half-wave plate (HWP) used to generate the linearly polarised light.

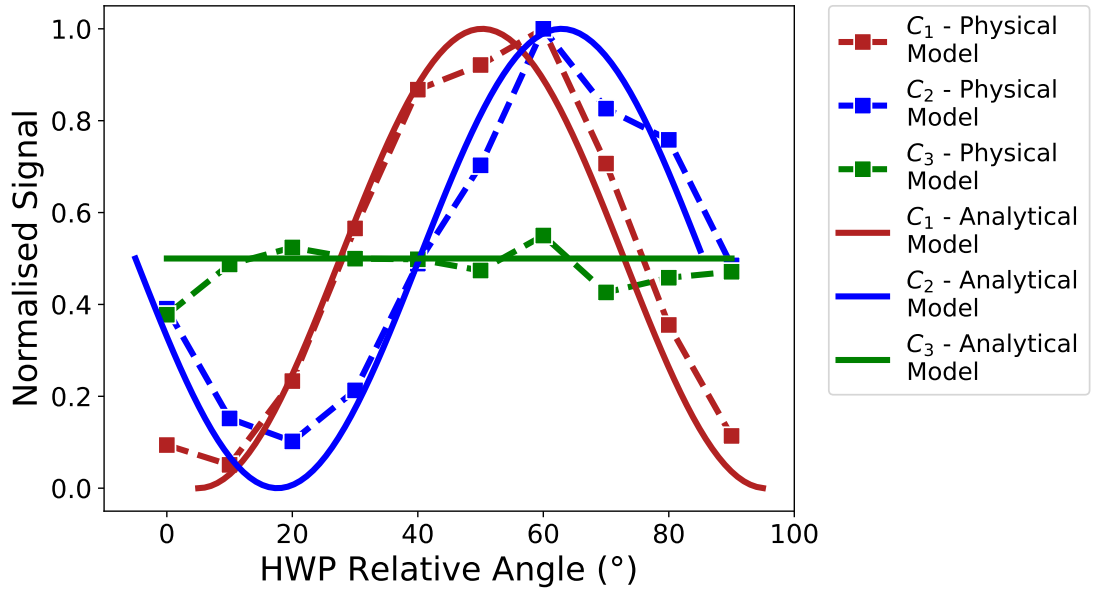


Figure 7.13: Comparison of the channel signals from the analytical and physical models of the polarimeter for variations in AOLP.

The resulting channel signals from the analytical model and the measurements made from the physical model are shown in figure 7.13. It is important to note that the x -axis in this case is labelled ‘relative HWP angle’, as the initial angle of the linearly polarised light was not known. Instead, starting from the initial AOLP, the HWP was rotated through 90° , corresponding to a 180° rotation of the AOLP.

Good agreement is observed between the two models, demonstrating that the polarimeter is capable of measuring changes in the AOLP.

The next variation examined was the degree of ellipticity in the polarisation of the light. In the analytical model, this was simple to achieve. For the physical model, a

quarter-wave plate (QWP) was placed between the source of linearly polarised light and the polarimeter. The QWP angle was varied incrementally from 0 to 90°, progressively increasing the degree of ellipticity and transitioning the light from linear to circular polarisation. A comparison between the two models is presented in figure 7.14.

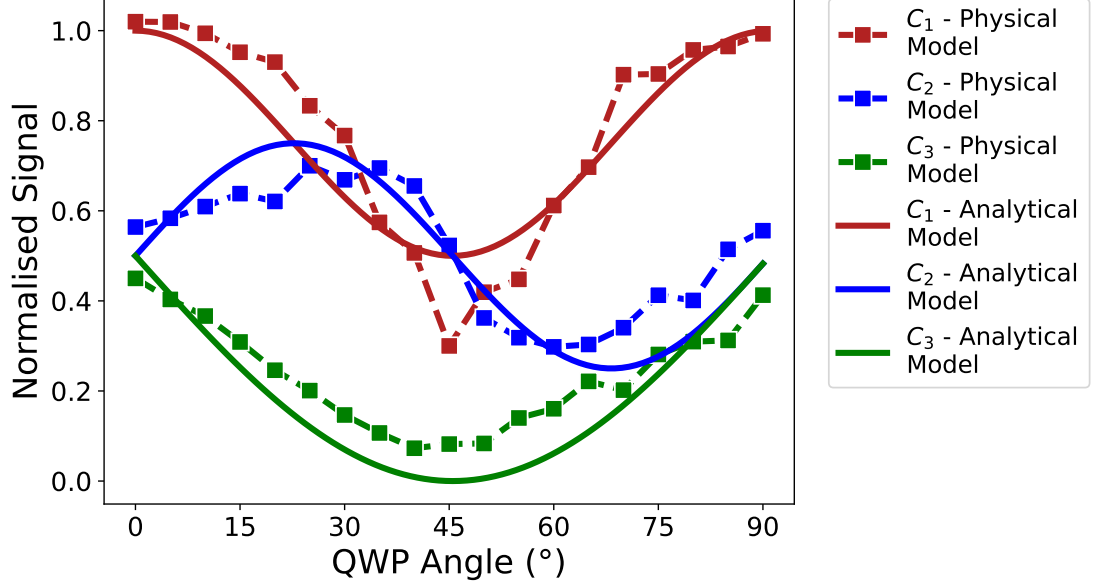


Figure 7.14: Comparison of the channel signals from the analytical and physical models of the polarimeter for light with varying ellipticity

Again, good agreement is observed between the two models. The results obtained from testing the operation of the polarimeter diagnostic give confidence in its correct functioning. This outcome shifted the focus toward exploring potential physical explanations for the experimental observations.

7.5 Comparing experimental data with the mathematical model of the polarimeter

Linearly polarised light

In an effort to find a physical explanation for the change in AOLP with transmitted light percentage, the channel signals measured experimentally during the pre-ioniser delay scan were compared with those from the analytical model. This comparison was based on the AOLP calculated from the experimental data. Initially, the assumption was made that no degree of ellipticity had been induced by the interaction, thus only

channels C_1 and C_2 (those used to calculate the AOLP) were compared. The results of this comparison are presented in figure 7.15.

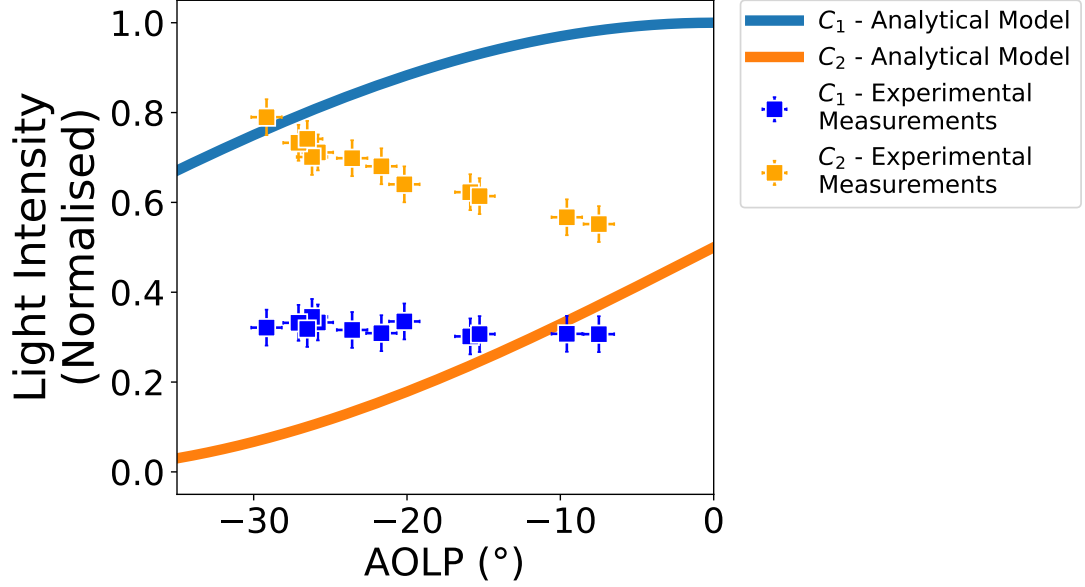


Figure 7.15: Channel signals produced by the analytical model of the polarimeter for varying input AOLP's. Also plotted are experimental measurements made by the polarimeter during the pre-ioniser delay scan shown in figure 7.5, plotted at their calculated AOLP.

The comparison shows a clear mismatch between the signal values predicted by the model and those measured experimentally. Under the assumption that the light entering the polarimeter is purely linearly polarised, such a discrepancy should not be possible. This leaves two further options, that the light entering the polarimeter may have been either unpolarised or possessed a degree of ellipticity, both of which are explored below.

Degree of polarisation

The next comparison between experimental and model values considered the case of light that was partially polarised. Physically, light with a degree of polarisation being generated from the interactions being studied would be challenging to explain. However, the comparison was carried out to definitively rule out this possibility. An experimental measurement from the pre-ioniser delay scan was selected, and within the analytical model the degree of polarisation was varied from unpolarised to polarised. While the degree of polarisation is varied, the AOLP remains unchanged. As a result, the corresponding channel values are represented by horizontal lines rather than discrete

data points, as seen in previous figures. If the degree of polarisation were responsible for the experimental measurements made by the polarimeter, the respective channel values would, at some point as the light changes from unpolarised to polarised, align with the analytical model values. This comparison is illustrated in figure 7.16.

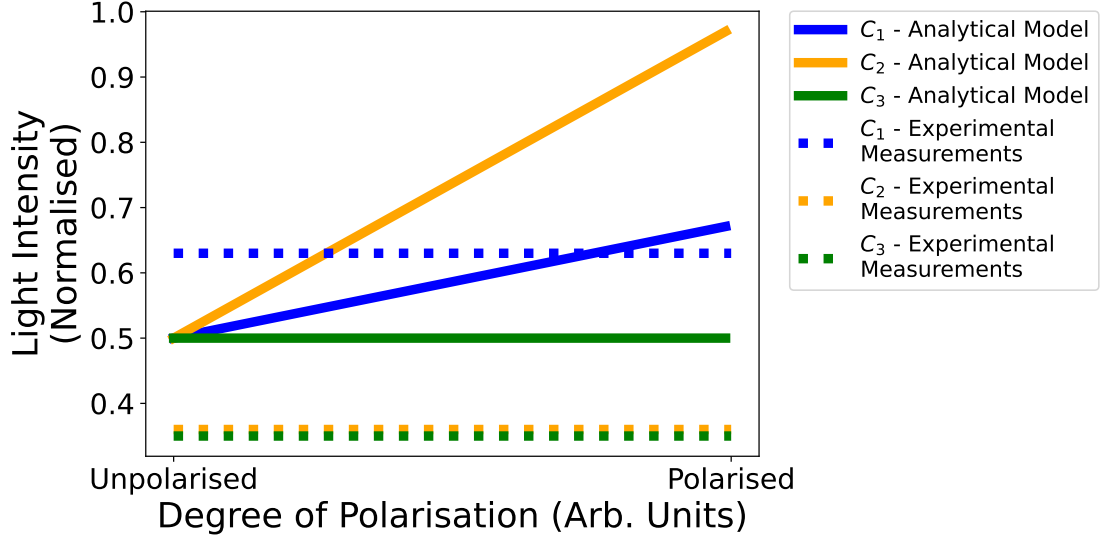


Figure 7.16: Channel signals produced by the analytical model of the polarimeter for varying degrees of polarisation. Also plotted are experimental measurements made by the polarimeter.

As there is no point at which the experimental measurements and the values taken from the analytical model align, the degree of polarisation can be ruled out as a cause of the unexpected results. This comparison was repeated for the remainder of the measurements made during the pre-ioniser delay scan and a lack of agreement was observed in all cases.

Elliptically polarised light

The final comparison was conducted for light with varying degrees of polarisation ellipticity. Similar to the investigation of the degree of polarisation, an experimental measurement from the pre-ioniser delay scan was selected, and within the analytical model the AOLP was set equal to the value derived from this measurement. The ellipticity was then varied incrementally, altering the polarisation state from linearly to circularly polarised. Figure 7.17 shows this comparison.

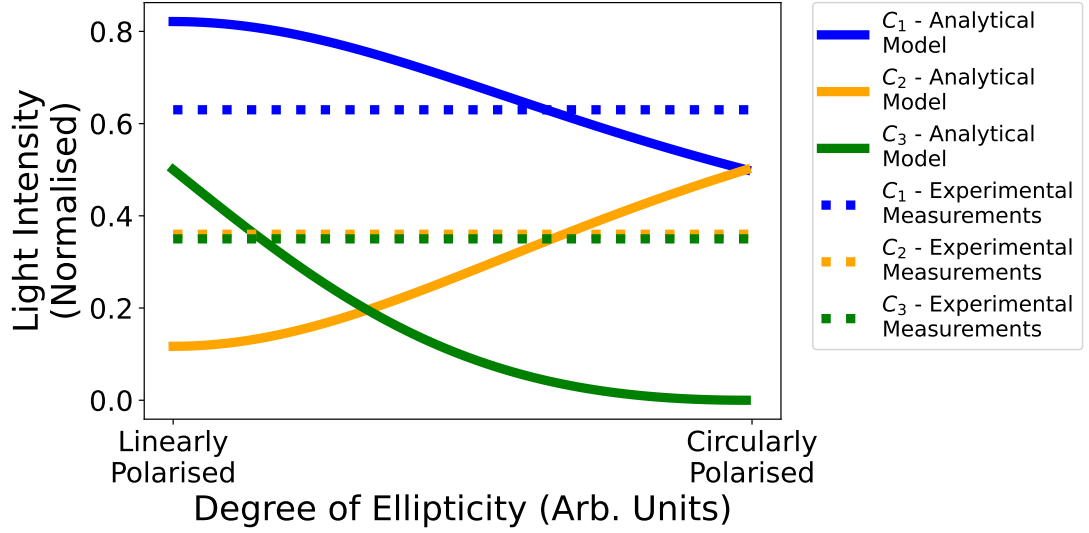


Figure 7.17: Channel signals produced by the analytical model of the polarimeter for varying degrees of ellipticity. Also plotted are experimental measurements made by the polarimeter.

As there is no point at which the experimental measurements and the values taken from the analytical model align, a degree of ellipticity can be ruled out as a cause of the unexpected results. This comparison was repeated for the remainder of the measurements made during the pre-ioniser delay scan and a lack of agreement was observed in all cases.

The results of the comparisons between the experimental and model values for the three potential polarisation states reveal that under none of the examined conditions do the values align. This indicates that the measurements obtained from the polarimeter during the experiment do not correspond to any physically possible single polarisation state. This indicates the need for a different approach to providing a physical explanation for the experimental results.

7.5.1 Effect of sub-sampling light for the polarimeter

Another potential reason for the polarimeter measurements is the way in which the light from the rear of the target was sampled. The design of polarimetry diagnostics used in prior experimental campaigns involved capturing all of the light from the rear of the target using a collection optic and re-imaging it to a series of CCD cameras, while during the experimental campaign being reported on in this chapter, only a sub-section of the light was collected. Shown in figure 7.18 are four examples of the transmission screen images captured during the experiment. From these images it can be seen that

only a small percentage of the light being investigated is passed into the polarimeter, with the spatial intensity of this light seeming to vary between shots.

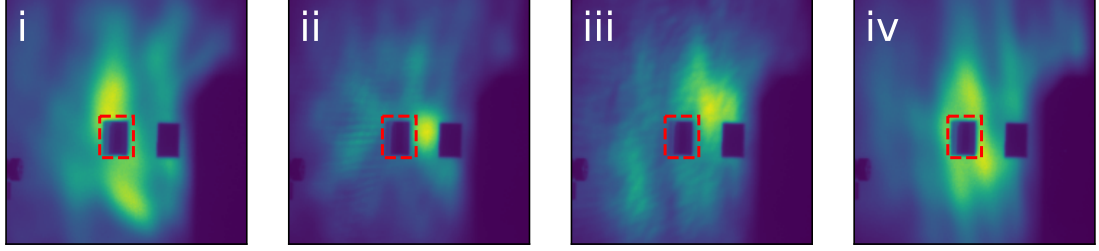


Figure 7.18: Examples of transmission screen images captured on four different shots during the experiment. The hole through which light was sampled for the polarimeter is highlighted in red.

If the polarisation state of the light varied both spatially and temporally during its generation, it is plausible that light with significant fluctuations in both intensity and polarisation state could enter the polarimeter. A possible source of these fluctuations is discussed in M.King *et al.* [58] in which the propagation direction of electrons being accelerated from the rear of targets undergoing transparency is shown to become increasingly unstable with increasing pulse duration, studied up to 0.6 ps. Given that electrons accelerated from the rear of the target are responsible for driving harmonic light emission, the use of 1 ps pulses in this experiment could have caused significant instabilities in the target as it became transparent. This could then lead to variation in the emitted light's properties that if present in the sub-sampled beam, could explain the measurements made during the experiment.

7.6 Modelling the interaction with 3D EPOCH simulations

To investigate the possibility that the transmitted and generated light at the rear of the target could have an intensity and polarisation state that vary both spatially and temporally, and if the sub-sampling of this light could result in the polarimeter signals measured experimentally, 3D EPOCH simulations were performed and analysed. For these simulations, the XYZ grid was defined as $1000 \times 720 \times 720$ mesh cells, corresponding to a simulation box volume of $20 \times 20 \times 20 \mu\text{m}^3$. A linearly polarised pulse with a wavelength of $\lambda_l = 1053 \text{ nm}$ was used, entering the simulation box from the left X boundary. The focal plane of the laser pulse was positioned $5 \mu\text{m}$ from the X boundary that the pulse entered from, with this plane defined as $X = 0$. All boundaries were

configured as free space. Both the spatial and temporal intensity profiles were defined as Gaussian, with FWHM of $\phi_L = 3 \mu\text{m}$ and $t_L = 400 \text{ fs}$ respectively. While this pulse duration is shorter than the 1 ps duration of the main pulse used during the experiment, the generation of light from the rear of the target necessary to test the hypothesis in question does still occur. This choice was made to reduce the computational resources required to perform these simulations. The peak pulse intensity was set to $I_L = 5 \times 10^{19} \text{ W/cm}^2$ and the simulation was performed for a total duration of 800 fs . The target density was initialised as Gaussian in space to mimic the generation of plasma caused by the interaction of a B8 pulse with the target prior to the arrival of the main pulse, with this Gaussian having a FWHM width of 30 nm . The target was comprised of C^{6+} and H^+ ions in a ratio of 1:6, neutralised by a population of electrons with a peak density of $30n_{\text{crit}}$. The data outputted from the simulations was comprised of E_X and E_Y field strengths, from which the polarisation could be calculated.

7.6.1 Spatio-temporal variation in light intensity and polarisation state

Prior to investigation of the simulation results, filtering the light for the ω_L frequency was performed as the polarimeter used during the experiment filtered for the fundamental frequency. This was achieved using a similar FFT function as was used in the previous results chapters of this thesis, covered in more detail in section 4.2.1.

The first aspect of the simulation results to be analysed was that of the spatio-temporal variations in the intensity and polarisation state of the light transmitted through and generated from the target. To do this, three line-outs through the simulation data were taken; one at $Z = 0$, through the center of the simulation box, another at $Z = 5 \mu\text{m}$ and the final at $Y = -5 \mu\text{m}$. The positions of the three line-outs are displayed in figure 7.19, labelled *A* to *C*. These are displayed on top of the time-integrated spatial profile of the light, measured $12 \mu\text{m}$ from the rear of the target, to aid in understanding their position relative to the light being examined.

The locations of the line-outs were chosen somewhat arbitrarily, as there was no prior knowledge of the existence or location of the properties of the light being investigated. As such, a reasonable approach was to select one line-out passing through the center and two additional off-center line-outs along different axes to increase the likelihood of varying polarisation states and intensities being discovered.

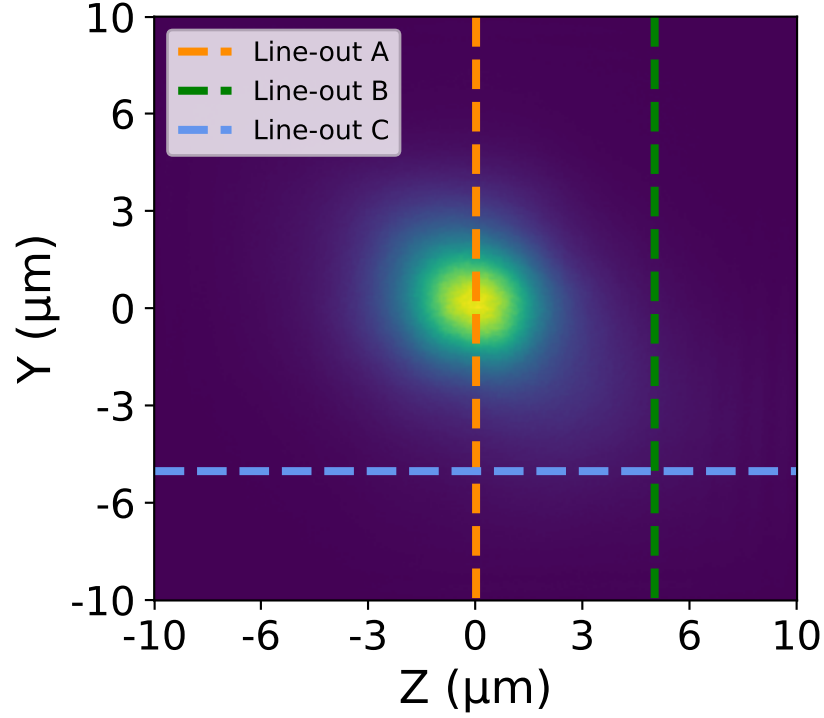


Figure 7.19: Positions of the three line-outs taken through the simulation data to measure spatio-temporal changes in polarisation state and intensity, relative to the time-integrated spatial profile of the light measured 12 μm from the rear of the target.

For each of these line-outs, both the AOLP and intensity of the light as it changes over time, measured 12 μm from the rear of the target, are plotted in figures 7.20. 7.21 and 7.22. Note that the x -axis in these plots has been truncated such that it begins at a simulation time of $t = 440$ fs. This choice was made as prior to this point, no light had been generated from or transmitted through the target. The method used to calculate the polarisation was to take the absolute values for E_Z and E_Y field strengths output from the simulation and calculate the resulting angle between the two vectors.

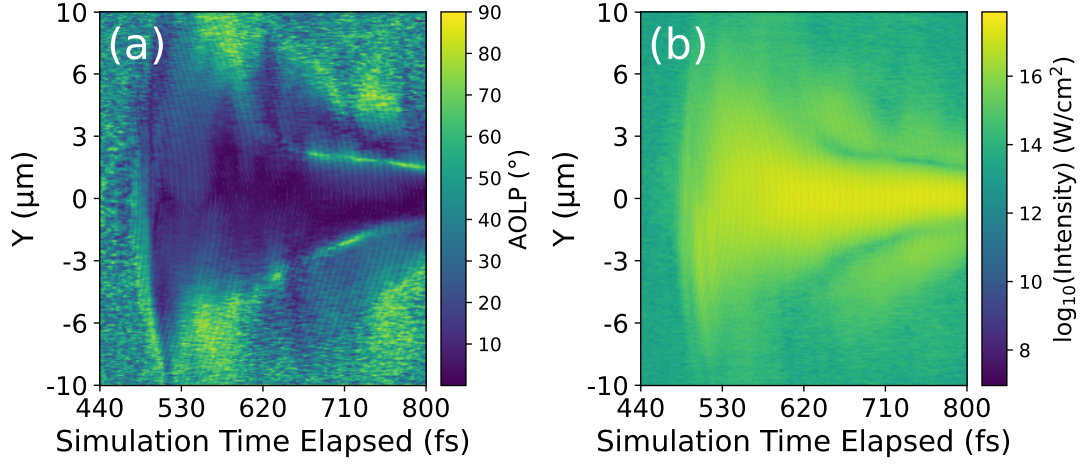


Figure 7.20: Change in (a) the polarisation state and (b) intensity over time and space for lineout A from figure 7.19. Results taken between a simulation time of $t = 440$ fs to the simulations completion.

In the case of line-out A, shown in figure 7.20(a), $\text{AOLP} = 0^\circ$ light is seen to be dominant near the laser propagation axis ($Y = 0 \mu\text{m}$) after $t \sim 650$ fs. This is likely due to the target becoming relativistically transparent to the main pulse at this time in the simulation, after which it is able to propagate through the under-dense plasma and the generation of E_Z is reduced. This light is also seen to be the most dominant in intensity terms, seen in figure 7.20(b), in which a more collimated beam is seen toward the end of the simulation. Prior to this point, between the simulation times of $t \sim 500$ and ~ 650 fs, the polarisation state can be seen to oscillate between 0 and $\sim 30^\circ$ in both space and time near the same axis. This may be attributable to the overlapping of E_Y and E_Z light generated from the rear of the target prior to the onset of transparency. Away from the laser propagation axis and at simulation times $\gtrsim 500$ fs, more significant changes in the polarisation state are observed (up to 90°), although with a much lower intensity. In addition to this, the spatial and temporal variations in this light appear to be more random.

In the case of line-out B, figure 7.21(a), oscillations with similar cyclical changes in polarisation state to those observed in figure 7.20(a) can be seen. Outside of this region, the changes seem essentially random. Figure 7.21(b) demonstrates that the oscillating component of the light has the highest intensity. Outside of this region, similar to figure 7.20(b), the light with the greatest degree of variation in polarisation state has a much lower intensity.

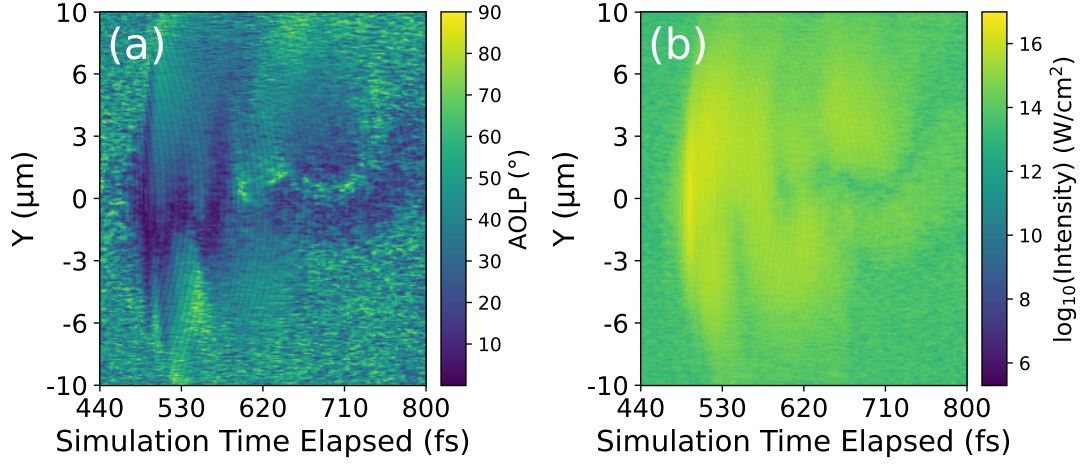


Figure 7.21: Change in (a) the polarisation state and (b) intensity over time and space for lineout B from figure 7.19. Results taken between a simulation time of $t = 440$ fs to the simulations completion.

Finally, the results from line-out *C* are shown in figure 7.22. Practically none of the oscillations present in the other two line-outs are present, with the variations appearing almost completely random. With respect to the intensity of the light, the areas with the greatest change in polarisation have the lowest intensity.

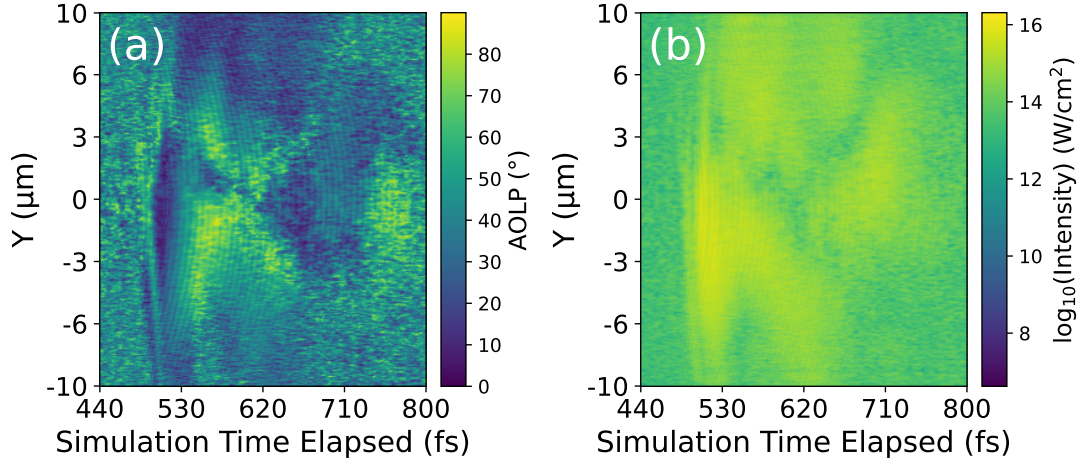


Figure 7.22: Change in (a) the polarisation state and (b) intensity over time and space for lineout C from figure 7.19. Results taken between a simulation time of $t = 440$ fs to the simulations completion.

Despite line-outs *B* and *C* being taken at equal distances from the center of the simulation box, though orthogonal to one another, there are significant differences observed in both intensity and polarisation state. A possible cause of this is the polarisation axis of the main laser pulse. As these two line-outs are taken along axes orthogonal to one another, the differences seen between them may be due to the dependence of light generation on the polarisation axis of the main pulse.

While the causes of the variation in the polarisation state of the light will be investigated further in a later section, the results displayed here support the hypothesis that the sub-sampling of the light by the polarimeter may explain the experimental results seen, warranting further analysis of the simulation data.

7.6.2 Recreating the sub-sampling of light from the rear of the target

The next stage of analysis was to attempt to recreate the way in which the light present at the rear of the target was sub-sampled for measurement by the polarimeter. One approach to this would be to first measure the size of the hole in the scatter screen used during the experiment. Then, with the distance from the target to the scatter screen known, this measurement could be used to calculate the corresponding size of the sampling region within the simulation. However, this method is complicated by the expansion of the light at the rear of the target during the experiment, which introduces additional uncertainties. As seen in the examples of images captured of the scatter screen during the experiment, shown in figure 7.18, it is likely that the beam overfills the scatter screen, making the necessary measurement of the total beam size not possible.

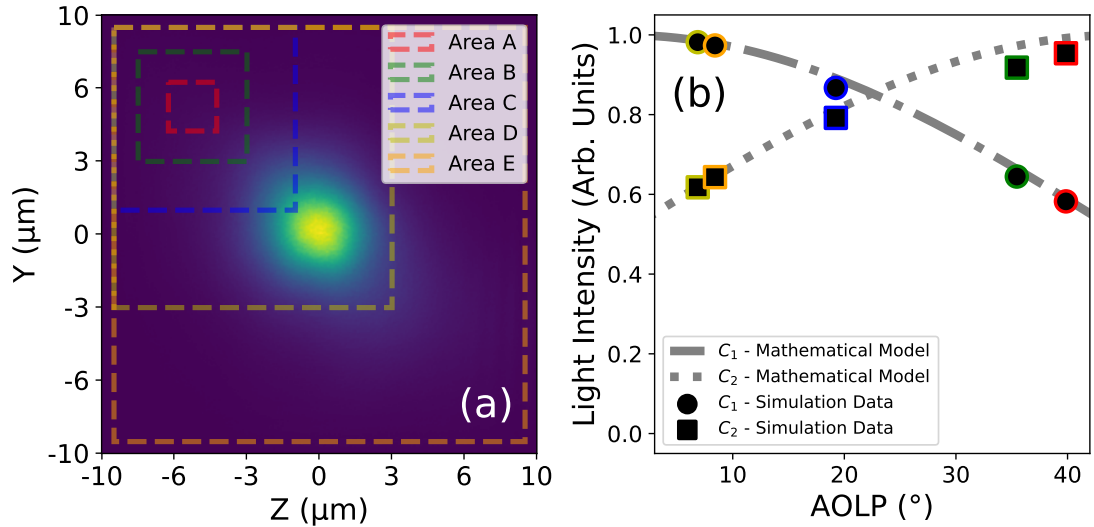


Figure 7.23: (a) The five sample size areas in which the polarisation of the light from the rear of the target was measured over the course of the interaction, overlaid onto the time-integrated spatial profile of the light measured 12 μm from the rear of the target. (b) The polarimeter channel values that resulted from measuring the polarisation state present in the five areas shown in (a), with the edge colour of the simulation measurements indicating the which of the sample areas they correspond to.

Instead, the choice was made to sample five areas of increasing size from the simulation data, shown in figure 7.23(a), overlaid onto the time-integrated spatial profile of the light, measured $12\text{ }\mu\text{m}$ from the rear of the target, illustrating their position relative to the light from the rear of the target. The approach was then to measure the polarisation of the light over the course of the interaction for each of the various sample areas. For each area illustrated in figure 7.23(a), the polarisation state at each time step was calculated and a Stokes vector representing both the polarisation state and intensity of the light was generated. This Stokes vector was then passed through a ‘virtual’ polarimeter, constructed using the py-pol package in a similar manner to the analytical model of the polarimeter described in section 7.4.1. This enabled the measurement of individual channel values on the polarimeter for all time steps in the simulations. To replicate the time-integrated nature of the polarimeter as used during the experiment, the channel values from each time step were summed, and an AOLP was calculated from these final channel values. The analytical model of the polarimeter was then employed to compare the expected channel values for the calculated AOLP with those obtained from the simulation. The comparison of simulation and model channel values for the five different area sizes is displayed in figure 7.23(b), with the edge colour of the simulation measurements indicating which of the sample areas they correspond to.

The results demonstrate minor discrepancies between the simulation and analytical model values for the smallest sample areas (areas A, B, and C). However, these discrepancies are not of the same magnitude as those observed in the experimental measurements. As the sample area increases in size, capturing the majority of the light emitted from the rear of the target, the simulation and model values are seen to align.

The analysis in this section assumes that the properties of light measured $12\text{ }\mu\text{m}$ from the target rear in the simulations are representative of those of the light when sampled at the scatter screen in the experiment, located tens of centimetres from the target. To support the validity of this assumption, the simulated field was numerically propagated by $40\lambda_L$ using the method described in [115]. The resulting spatial intensity profile is shown in figure 7.24.

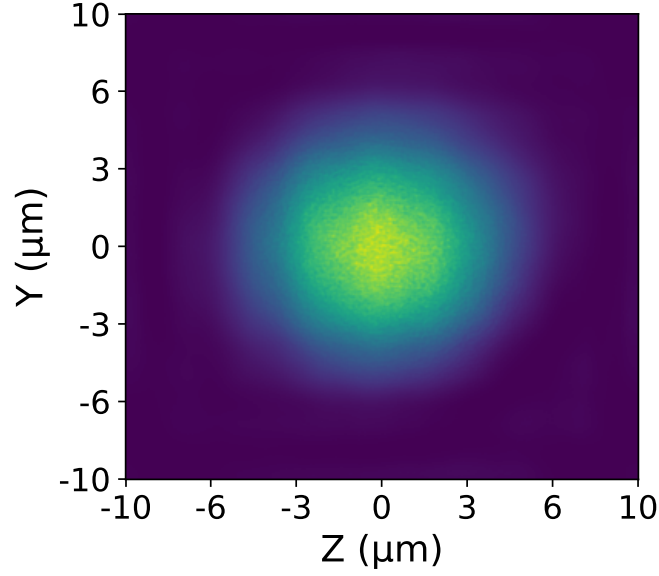


Figure 7.24: Spatial intensity profile of the simulated light after numerical propagation by $40 \lambda_L$ from the sampling plane located $12 \mu\text{m}$ behind the target rear.

Although this propagation distance is far shorter than the experimental target-screen separation, simulation size and resolution limit the achievable propagation range. Nevertheless, the similarity between the spatial profiles in figures 7.23(a) and 7.24 suggests that near-field sampling in the simulations provides a reasonable approximation of the far-field light distribution observed at the scatter screen.

Introducing sample box drift

When comparing the spatial profile of the light at the rear of the target from the simulation to that obtained from the transmission diagnostic images, the latter displays a significant degree of randomness not present in the former. This can be seen when comparing the images in figure 7.18 captured during the experiment to the time-integrated spatial profile from the simulations shown in figure 7.23(a). This suggests the presence of movement in the light as it is generated and transmitted by the target during the experiment that is absent from the simulations. In an experimental setting, a multitude of additional factors, such as minor variations in thickness within a single target, variations in pre-pulse intensity and uncorrected wavefront aberrations, have the potential to introduce asymmetries of a much greater magnitude than those observed in simulations. If, during the experimentally investigated interaction, the light from the rear of the target exhibits a sufficient degree of movement as it is generated and as

transparency occurs, this could contribute to a greater spatio-temporal variability in the polarisation state and intensity that is observed in the experimental measurements.

To account for this increased variability, the smallest sample area from figure 7.23(a) was selected, and a random ‘drift’ in its position was introduced over the course of the simulation, while again measuring the polarisation state. The movement of the sample area throughout the duration of interaction is depicted in figure 7.25, overlaid on the time-integrated spatial profile of the light measured $12\text{ }\mu\text{m}$ from the rear of the target.

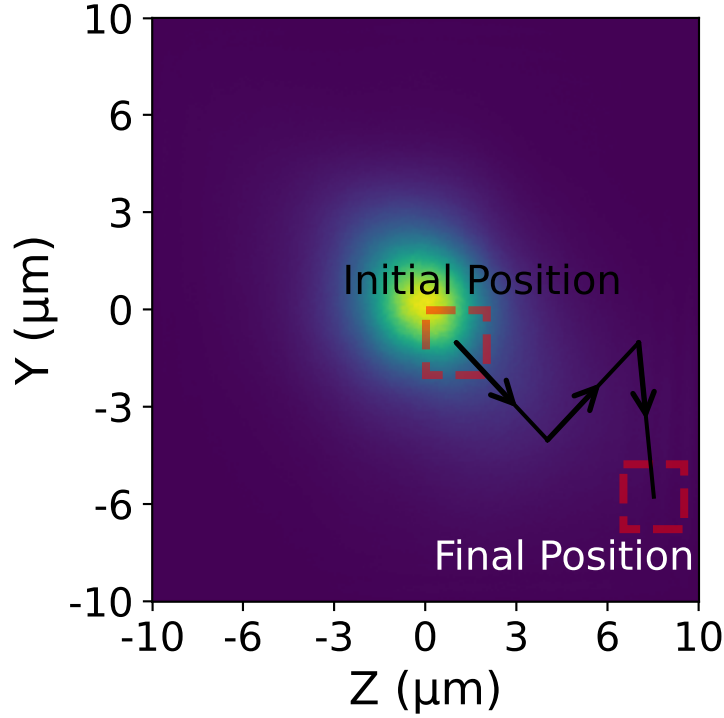


Figure 7.25: Illustration of the random movement of the sample area in which the polarisation was measured over the course of the interaction, overlaid onto the time-integrated spatial profile of the light measured $12\text{ }\mu\text{m}$ from the rear of the target.

Figure 7.26 shows the polarimeter channel measurements obtained from the simulation with the inclusion of sample box drift, compared to the analytical model. Additionally, the channel values for a selected experimental measurement are included. The introduction of drift results in channel values that deviate from the model values with a magnitude comparable to those observed in the experimental measurements. This suggests that the measurements made by the polarimeter during the experiment are explainable by the sampling of only a portion of the light emitted from the rear of the target, combined with spatio-temporal variations in the polarisation state and intensity of this light over the course of the interaction.

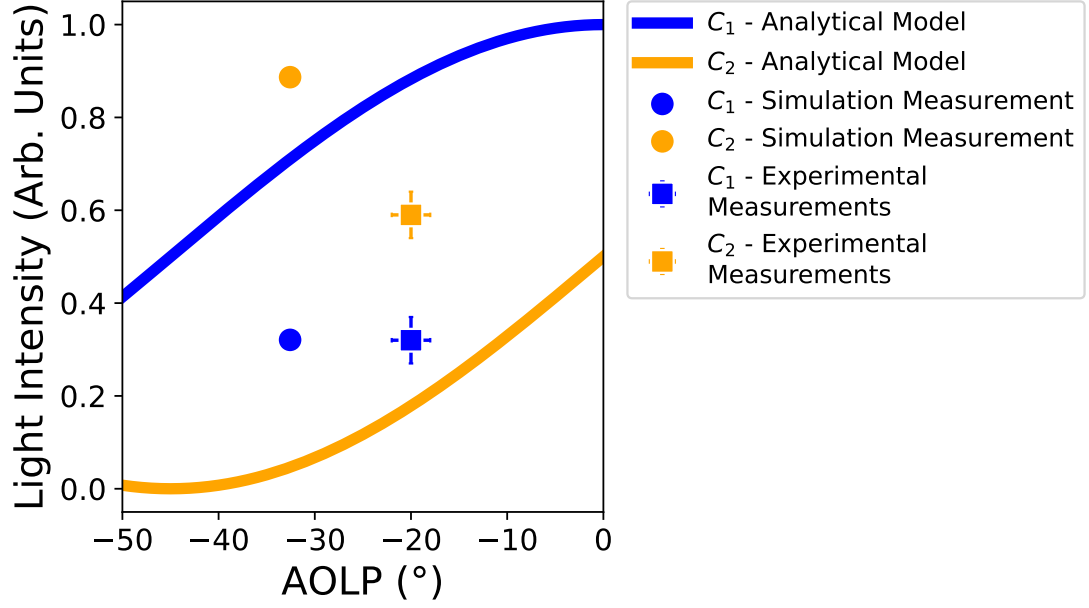


Figure 7.26: Channel values that resulted from a measurement of the polarisation state in the area illustrated in figure 7.25 in comparison with the channel values produced by the analytical model of the polarimeter. Also included are experimental measurements made by the polarimeter from one shot in the pre-ioniser delay scan.

7.6.3 Source of polarisation variation

The previous section establishes that a potential cause of the polarimeter measurements made during the experiment is the degree of variability in both the intensity and polarisation state of the light at the rear of the target, both spatially and temporally. In an attempt to understand the cause of this variability, the electron dynamics within the target were examined over the course of the simulation. This began by investigating changes in the electron energy density over time in four areas of the target. Energy density was chosen as a metric to examine as the electrons drive harmonic light emission at the target rear and therefore investigation of the electron dynamics could provide additional insight into the generation of this light. The metric of electron density alone was not chosen as this would not necessarily have represented the electrons that are leaving the target to generate light. Energy density in this context is found through the product of the electron number density and the corresponding electron energy as output from the simulation for a given time-step. The four areas of the target selected are displayed in figure 7.27(a).

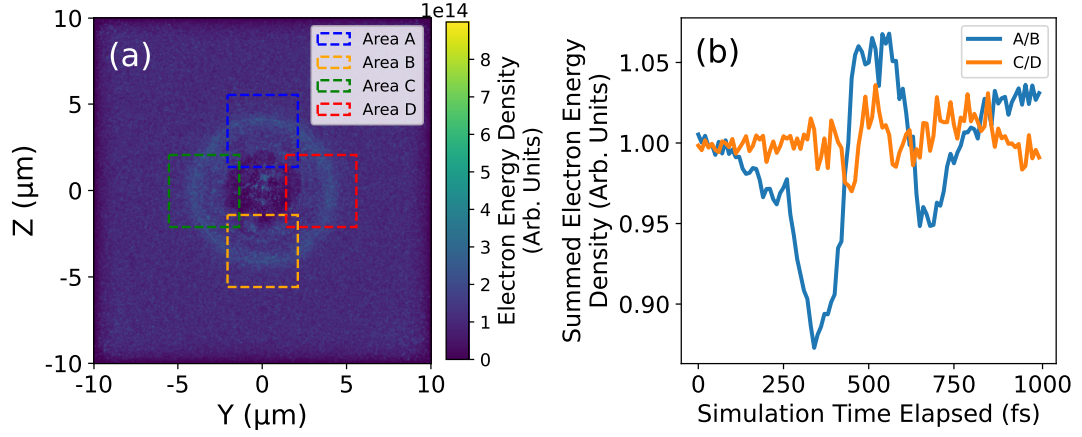


Figure 7.27: (a) The four sample areas in which the electron energy density was measured over the duration of the simulation, overlaid onto the target electron energy density profile present at $t = 740$ fs. (b) The change in the ratio of the summed electron energy densities in areas A/B and C/D over the duration of the simulation.

To analyse the temporal changes, ratios of the summed electron energy densities between opposite sample areas were calculated, those being areas A/B and areas C/D. This methodology was chosen as it was thought that any deviation of these ratios from a value of 1 would be evidence of an asymmetry in the acceleration of electrons from the target and the subsequent generation of light that contributes to spatio-temporal variations in polarisation and intensity. The variation of the two ratios over the course of the simulation are presented in figure 7.27(b). For the C/D ratio, measured along the non-polarisation axis, the variation is minimal, deviating from 1 by only a few percent. In contrast, for the A/B ratio, taken along the polarisation axis, a stronger asymmetry is observed. A maximum deviation from 1 of 12% is observed. It is possible that this asymmetry in electron energy density indicated by the variation in the A/B ratio is responsible for the spatio-temporal changes in polarisation and intensity.

Figure 7.28(a) shows the A/B ratio over the course of the interaction (labelled ‘Simulation A’), while figures 7.28(b) and (c) illustrate the changes in the intensity of the E_X and E_Y fields, respectively, measured $12 \mu\text{m}$ from the rear of the target. When examining this ratio, it is evident that the period of maximum deviation from 1 roughly corresponds to the time frame during which the greatest spatial variation in field intensities occurs. As the polarisation shifts observed are a result of the overlap of these two fields, this provides evidence supporting the conclusion that the spatio-temporal variation in the polarisation state can be attributed to the observed asymmetry in the electron energy density.

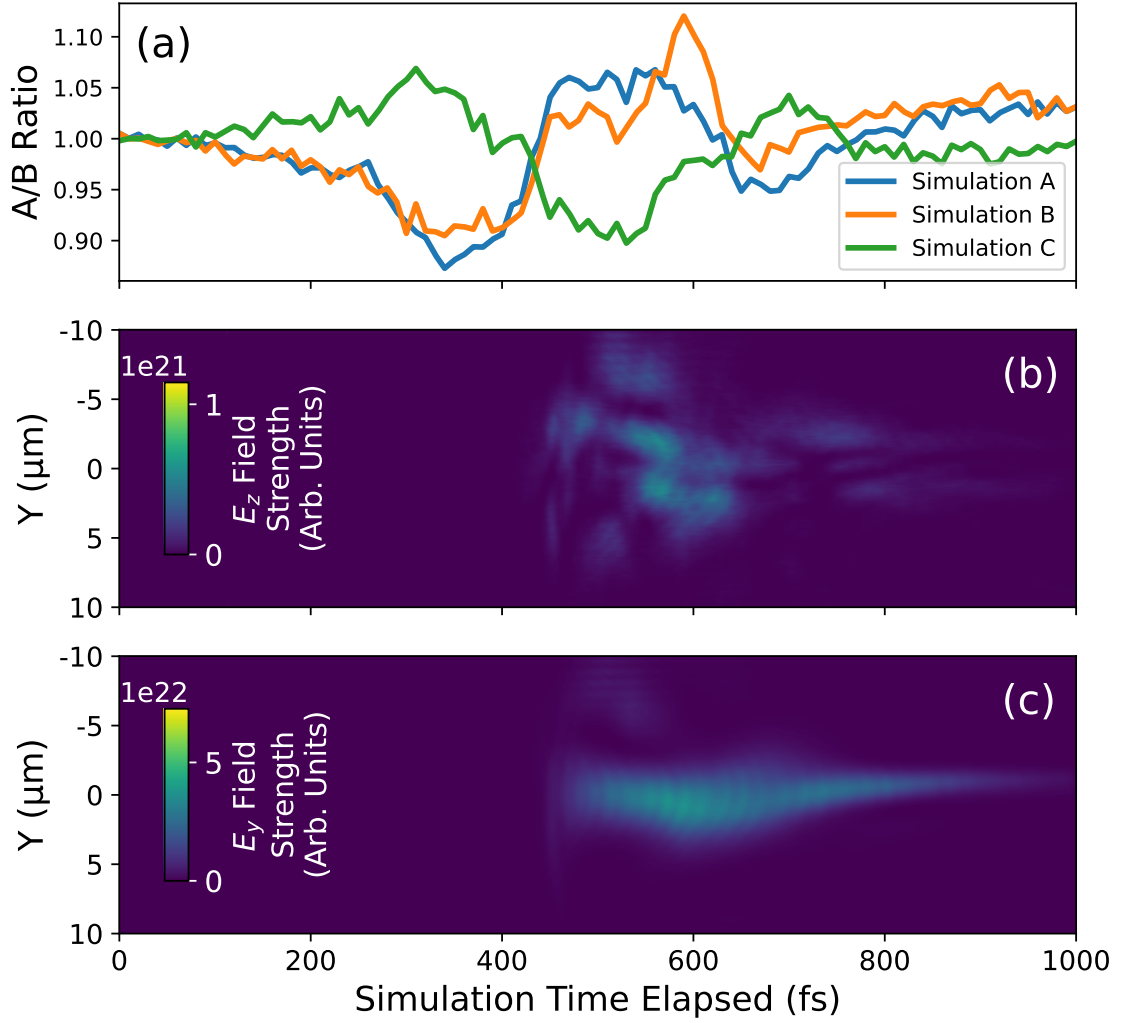


Figure 7.28: (a) The change in the ratio of the summed electron energy densities in areas A and B as shown in figure 7.27 for over the duration of the three simulations performed. (b) and (c) the change in the E_x and E_y fields respectively over the course of the interaction measured $12 \mu\text{m}$ from the rear of the target, taken from Simulation A.

In an effort to understand the sensitivity of the asymmetry to initial interaction conditions, and thus potentially its cause, two additional 3D simulations were performed. In both simulations, the initial particle distribution was randomised to a small degree relative to the first simulation, and in the second of the additional simulations, the phase of the laser as it first interacts with the target was changed by π . The same ratio of electron energy densities in sample areas A and B are plotted in figure 7.28(a), labelled ‘Simulation B’ and ‘Simulation C’ respectively.

When comparing the electron energy densities from Simulation B to Simulation A, a notable difference is seen. The maximum deviation from a ratio of 1 is positive and occurs at a different time in the simulation, suggesting that the electron dynamics, and

thus the light generation, is sensitive to minor changes in the initial electron distribution. When examining this metric for Simulation C, an increase in the ratio is observed at a simulation time of $t \sim 150$ fs, whereas both Simulations A and B exhibit a decrease at this point. Conversely, between simulation times of $t \sim 400$ fs and 600 fs, Simulation C shows a decrease, while the other simulations display the opposite behaviour. This pattern suggests that altering the phase of the laser results in a shift in the side of the target from which electrons are preferentially accelerated at a given time, relative to Simulations A and B. This demonstrates a sensitivity to the phase of the laser as it begins interacting with the target.

In the context of simulations, a minor alteration to the particle distribution within the target or a change in the laser phase results in relatively small changes to the electron dynamics. However, in an experimental setting, a multitude of additional factors, such as variations in the pulse contrast and uncorrected wavefront aberrations, have the potential to introduce asymmetries of a much greater magnitude than those observed in simulations. Such asymmetries could be responsible for the greater variability in spatial profile of the light at the rear of the target that was measured experimentally, which are likely be accompanied by larger spatio-temporal changes in intensity and polarisation state that may be the cause of the polarimeter measurements made during the experiment when only sampling a relatively small region of the transmitted/generated beam.

7.7 Conclusion

This chapter reports on the results of an experimental campaign conducted to investigate the changes in the polarisation state of light transmitted through ultra-thin (tens of nm) targets. This built on previous work in which a change in AOLP was found to be dependent on the percentage of the main pulse that was transmitted through the target. This change in polarisation is hypothesised to have occurred due to the super-position of light fields from the rear of the target and the light that propagated through the target after the onset of relativistic transparency, previously observed experimentally for pulse durations on the order of tens of fs. The experimental campaign reported on in this chapter was conducted using the Vulcan laser system, generating picosecond duration pulses. The aim was then to determine if similar apparent rotations

in polarisation state can be observed when such longer pulse durations are used.

The initial stage of the experiment involved an investigation into how the percentage of transmitted light varied with target thickness, ranging from $l = 10$ nm to $l = 125$ nm. Once 25 nm targets were identified as optimal for achieving a suitable transmission percentage, a second, lower-energy pulse was introduced as a pre-ioniser. By varying the delay between the pre-ioniser and the main pulse, the aim was to control the degree of target expansion and, consequently, the transmission percentage and AOLP. Although the expected trends in transmission percentage with pulse delay were not observed, the variations in transmitted light percentage were deemed sufficient for studying AOLP changes. As a novel polarimetry diagnostic was employed during this experiment, the efforts made to verify the method used to calculate the AOLP are covered due to the central role played by the AOLP measurements in the objectives of the experimental campaign. Unexpectedly, the relationship between transmitted light percentage and AOLP appeared random, contrary to prior studies. This shifted the focus to explaining these results.

To ensure that the polarimeter was functioning correctly, two models of the diagnostic were constructed, a physical replica and an analytical model. Comparing the responses of both models to light with varying polarisation states demonstrated agreement, confirming the polarimeter was functioning as expected and able to measure AOLP changes accurately. With the diagnostic validated, experimental measurements were compared to an analytical model of the polarimeter. This comparison suggested that the measurements were not consistent with any single polarisation state, leading to the hypothesis that the sub-sampling of light from the rear of the target, combined with spatio-temporal variations in light intensity and polarisation state, contributed to the unexpected results. This may be an explanation as to why the experiment reported on in M. Duff *et al.*[56], in which the polarimeter employed sampled the entire beam, did not report experimental measurements similar to those discussed in this chapter.

To test this hypothesis, 3D EPOCH simulations were conducted. These simulations revealed spatio-temporal variations in polarisation and intensity, and when sub-sampling was applied to the simulated data, results resembling the experimental measurements were obtained. Finally, an exploration of the underlying cause of the polarisation variations was conducted by analysing the electron dynamics within the target during the interaction. This analysis found asymmetries in the electron energy

densities over the course of this simulation which could be a contributing factor to the spatio-temporal variations in polarisation and intensity. By performing additional simulations, these asymmetries were found to be sensitive to minor changes in the initial interaction conditions such as particle distribution and laser phase.

This study identifies spatio-temporal variations in both the polarisation and intensity of light generated and transmitted through ultra-thin targets undergoing RSIT due to interactions with relativistically intense pulses. The somewhat random nature of these variations and the sub-interaction times scales on which they occur highlight the need to measure the majority of the light generated from such interactions when measuring its polarisation state, which may inform the design of future polarimeter diagnostics.

Chapter 8

Conclusions & future work

The overarching objective of the work presented in this thesis is the investigation of laser-solid interaction based relativistic plasma photonics techniques for the purposes of polarisation control and higher-order mode generation. This has been achieved through numerical and experimental studies, beginning with a simulation-based investigation of the generation of $2\omega_L$ light with higher-order spatial modes from the interaction between the relativistically intense laser pulses and pre-formed aperture targets. This was followed by a numerical study of experimental challenges, an experimental campaign on $2\omega_L$ generation, and finally an experimental investigation of polarisation control using ultra-thin targets. This chapter summarises the key findings of these investigations, outlines the main conclusions drawn, and briefly discusses potential directions for future research.

Chapter 4 reports on an investigation into the use of a preformed plasma aperture to generate intense second harmonic light with higher order modes, explored through 2D and 3D EPOCH simulations. This builds on the work previously published by M. Duff *et al* [56], in which light with similar properties was generated from relativistic plasma apertures formed during the interaction of intense laser light with an ultra-thin targets. It is demonstrated that a laser pulse with a Gaussian spatial profile incident on a preformed aperture target can produce relativistically intense second-harmonic light with higher-order spatial modes with conversion efficiencies up to 5% in 2D and 3.5% in 3D simulations. For a linearly polarised pulse with $\phi_L = 2.1 \mu\text{m}$ and $I_L = 1 \times 10^{21} \text{ W/cm}^2$, these figures were seen for aperture diameters of 2 and 3 μm respectively, with the conversion efficiency reaching a maximum and subsequently saturating for target thickness of $\sim 7 \mu\text{m}$ and greater. At higher intensities, relativistic transparency at the

aperture edges reduced efficiency by effectively enlarging the aperture diameter. The use of 3D simulations allowed for the use of circularly polarised light, demonstrating broadly similar trends in conversion efficiency.

Analytical modelling demonstrated that efficiency scales with the strength of the longitudinal electric field at the aperture edge, the fraction of incident laser energy absorbed by the target, and the proportion of the incoming light transmitted through the aperture. The lower efficiency observed in 3D was attributed to reduced interaction between fields and aperture edges when the polarisation vector is misaligned, and to additional degrees of freedom allowing electron escape.

An inverse relationship between the aperture diameter and the size of the generated mode structure was observed, demonstrating the ability to tune the mode size through varying target geometry. The influence of the presence of plasma within the aperture channel was studied, with densities ranging from $0.01n_{crit}$ to n_{crit} investigated. The presence of plasma within the aperture channel was found to reduce efficiency across all densities studied, with stronger reductions at higher densities. Finally, tapered apertures matched to the convergence angle of the self-focusing pulse improved performance, achieving efficiencies up to 6.9%.

In chapter 5, the potential challenges associated with the experimental investigation of $2\omega_L$ generation from aperture targets were explored using numerical simulations, with the aim of developing strategies to mitigate these effects. Two key factors were identified as potentially influencing the interaction dynamics. The first is the formation of pre-plasma within the aperture channel along the laser propagation axis due to the ASE pedestal and pre-pulse. The second is the shot-to-shot spatial jitter of the laser pointing.

EPOCH simulations revealed that even modest pre-plasma expansion reduces $2\omega_L$ yield, motivating the use of FLASH hydrodynamic modelling. FLASH simulations showed aperture narrowing at lower pre-pulse intensities and near-complete filling at higher intensities. This justified the use of combined FLASH and EPOCH simulations to investigate pre-plasma effects on both $2\omega_L$ generation and mode structure. The results demonstrated that, under the planned experimental conditions, the desired spatial mode would still be generated with an acceptable yield. Further combined FLASH and EPOCH simulations examined whether the presence of pre-plasma alters the previously identified optimal focal spot to aperture diameter ratio. The results indicated a

reduction in the optimal ratio from ≈ 1 to the range 0.7-0.9. Given the planned target diameters ($d = 6$ and $10\ \mu\text{m}$) and the achievable focal spot size at best-focus ($3\ \mu\text{m}$), this ratio could not be realised under the planned experimental setup. A possible solution was identified in which the laser is deliberately defocused to enlarge the focal spot at the point of interaction, thereby achieving the optimal ratio at the cost of reduced intensity.

The influence of spatial jitter was also studied using FLASH-EPOCH simulations, benchmarked against Gemini laser jitter measurements. Results showed that jitter can both reduce yield and distort mode structure, although in some cases a small degree of jitter enhanced performance by bringing high-intensity regions into contact with the aperture edge. Two strategies were identified to mitigate the negative effects of jitter. The first is a statistical approach, relying on a large number of experimental measurements to increase the probability of achieving favourable conditions. The second is to position the target away from best focus, creating a large focal spot at the point of interaction. This approach reduces the relative sensitivity to jitter while simultaneously improving the spot size to aperture diameter ratio, though again at the expense of reduced pulse intensity and, consequently, a lower $2\omega_L$ yield.

Chapter 6 reported on an experimental campaign using the Gemini laser to generate $2\omega_L$ light from aperture targets of 6 and $10\ \mu\text{m}$ diameter. Measurements of transmitted ω_L and generated $2\omega_L$ light were made as functions of pulse energy and focal spot size, with the spot varied by shifting the target relative to the focusing optic over three positions: $\Delta_X = 20$ and $29\ \mu\text{m}$.

Results were compared with FLASH-EPOCH simulations, showing good agreement when the focal spot was near-Gaussian and stable. Under such conditions, the spread in $2\omega_L$ yield could largely be attributed to spatial jitter. However, when the spot was enlarged by defocusing, deviations arose that could not be explained by jitter alone. Instead, non-Gaussian spot profiles introduced by defocusing were identified as the dominant cause.

The findings highlight that experimental studies of aperture targets are highly sensitive to both the magnitude of spatial jitter and the quality of the focal spot spatial intensity profile. While defocusing can, in principle, mitigate jitter and improve spot-to-aperture matching, this benefit is undermined if the spot profile degrades significantly. Future experiments may instead benefit from optics with higher F-numbers, producing

larger Gaussian focal spots at best focus and reducing reliance on defocus strategies.

Chapter 7 presents the results of an experimental campaign investigating changes in the polarisation state of light transmitted through ultra-thin (tens of nm) targets. This builds on previous work [56] which reports changes in the angle of linear polarisation (AOLP) of light transmitted through such targets. This is attributed to the superposition of light emitted from within the target with the light transmitted through it after the onset of relativistic self-induced transparency (RSIT), observed experimentally with pulse durations of tens of femtoseconds. The present work, conducted using the Vulcan laser system, extended this investigation to the use of picosecond-duration pulses, with the aim of determining whether similar polarisation changes could be observed with the use of longer duration pulses.

A pre-ioniser pulse was used to control target expansion and hence the transmission of the main pulse. Although expected trends in transmission as a function of delay were not observed, sufficient variation was achieved to study the angle of linear polarisation (AOLP). Measurements revealed an apparently random relationship between transmission and AOLP, inconsistent with simple models and requiring further efforts to understand these results. Diagnostic tests confirmed the polarimeter’s accuracy, suggesting that the observed behaviour originated in the interaction itself. This may also account for the absence of similar observations in the experiment of M. Duff *et al.* [56], where the diagnostic sampled the entire beam rather than a sub-section.

Comparisons with analytical models and 3D EPOCH simulations indicated that sub-sampling of the transmitted beam, combined with spatio-temporal variations in intensity and polarisation, can explain the experimental observations. These variations were traced to asymmetries in electron energy densities within the target, which were highly sensitive to initial conditions. The study therefore identifies intrinsic spatio-temporal fluctuations in polarisation during relativistic transparency, emphasising that accurate characterisation requires capturing the majority of transmitted light. This conclusion provides important guidance for the design of future polarimetry diagnostics in high-intensity laser-plasma experiments.

The investigations presented in this thesis demonstrate both the opportunities and the challenges of using relativistic plasma photonics to manipulate light in laser-solid interactions. Taken together, the results demonstrate that relativistic plasma photonics offers routes for controlling light at extreme intensities, but also depends critically on

precise control of target design, laser stability, and diagnostic methodology.

8.1 Future work

The numerical investigations presented in chapters 4 and 5, together with the experimental campaign described in chapter 6, lay the foundation for further exploration of relativistic plasma apertures as a method for generating frequency-doubled light with higher-order spatial modes. Many of the features of this mechanism discussed in chapter 4 have not yet been examined experimentally. While the production of frequency-doubled light has been demonstrated, direct imaging of the generated mode has not yet been achieved. This, along with the potential to tailor the mode shape through varying interaction parameters such as pulse polarisation and aperture diameter, represents a promising aspect of this plasma photonic mechanism that warrants further investigation.

Although only briefly addressed in chapter 4, another promising direction is the use of aperture targets with more complex geometries. This could include tapered aperture channels, as outlined in chapter 4, or alternative transverse geometries, such as those discussed in reference [63]. These provide potentially fruitful aspects of this plasma photonic mechanism that remain to be confirmed experimentally, and the experimental campaign described in chapter 6 has highlighted the sensitivity of such investigations to spatial jitter and focal-spot quality, suggesting that future experiments will benefit from improved pointing stability and larger, well-preserved Gaussian focal spots, potentially achieved with higher F-number optics.

Beyond aperture targets, the control of the polarisation state of light transmitted through ultra-thin targets undergoing RSIT for longer pulse durations (picoseconds) remains an open question. However, the findings of chapter 7 demonstrate that sub-sampling the transmitted light can introduce significant measurement inaccuracies. This motivates the development of new polarimetry diagnostics capable of capturing the full transmitted light, which will not only advance studies of relativistic plasma photonics but also improve the accuracy of polarisation measurements across a wide range of high-intensity laser-plasma experiments.

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