

In situ diagnostics and stabilisation of plasma-based particle accelerators towards applications

by

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Declaration

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Abstract

Plasma-based particle acceleration offers a novel pathway for achieving ultra-high energy gains in charged particles over propagation distances orders of magnitude shorter than those required by conventional RF accelerators. Demonstrated beam properties, including femtosecond bunch duration, low divergence, and high peak current, underscore its potential for deployment across a broad range of scientific and industrial use cases. Yet, plasma accelerators are presently limited to research-grade facilities, with their successful operation contingent upon the support of extensive, expert personnel.

This work examines three topics that each play a critical role in advancing plasma accelerators toward practical, application-ready machines. The first one is driving the accelerator with more accessible, kHz laser systems. We experimentally validate the e-KAIO (electron Kit-All-in-One) laser-plasma accelerator module for kHz laser operation. This constitutes a first step towards the wide dissemination of fully integrated laser-plasma accelerator technology. Furthermore, we demonstrate that existing kHz laser-plasma accelerators are capable of generating electron fluxes with characteristics comparable to those observed in the Earth's Outer Van Allen Belt, fulfilling the rising demand for beams sources capable of carrying out accelerated lifetime radiation hardness testing on space-qualified electronics.

The second topic is the stabilisation of accelerated particle beams. Here, we show that employing fluctuating output electron beams from an LWFA to drive

an additional PWFA stage with a plasma photocathode injector, can produce secondary electron beams with greater stability, higher quality, and improved reliability. This hybrid approach also acts as a brightness transformer with 5D and 6D brightness values of $B_{5,n} > 10^{19} \text{ A m}^{-2} \text{ rad}^{-2}$ and $B_{6,n} > 10^{18} \text{ A m}^{-2} \text{ rad}^{-2} / 0.1\% \text{ BW}$ achieved during this study, nurturing prospects for demanding applications such as free-electron lasers.

The third theme is the establishment of a live non-intrusive accelerator diagnostic. Here, we construct a well-controlled experimental setting in which correlations between the plasma afterglow signal and key accelerator parameters: laser pointing angle, laser energy and gas density, can be quantitatively assessed, enabling the emitted light to be leveraged as a diagnostic for quantities that have historically been difficult to measure at the interaction point. Additionally, the scope of the afterglow technique was broadened by illustrating its ability to diagnose the non-linear ionisation-induced defocusing experienced by an intense laser pulse propagating in a gaseous medium, distinguishing this method from comparable diagnostics.

Each part of this work advances its respective research domain, yet collectively it lays the foundation for generating electron beams with orders of magnitude higher brightness, strength and frequency in next-generation staged plasma accelerator architectures.

Acknowledgements

Reflecting on the past 5 years, I feel incredibly fortunate to have been able to dedicate a portion of my life to carrying out research. Most days felt both stimulating and rewarding, providing me the space to learn new skills and understand physics on a deeper level in the pursuit of developing new technologies. That being said, this degree has comfortably been the most challenging thing I have ever undertaken and has taught me how to continually persevere and carry on in the face of adversity. Now I'm on the other side of those late-night lab stints and weekend library sessions, I can say that I am forever grateful for this.

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Publications

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Table of Contents

1	Introduction	1
1.1	Particle accelerators	1
1.2	Thesis outline	3
1.3	Role of the author	4
2	Fundamental Concepts	6
2.1	Electrodynamics	6
2.2	Laser beams	9
2.3	Plasma physics	12
2.3.1	Plasma generation	13
2.3.2	Ionisation-induced defocusing	19
2.3.3	Plasma evolution	20
2.4	Laser propagation in a plasma	22
2.4.1	Relativistic self-focusing	23
2.4.2	Self-phase modulation	25
2.5	Plasma wakefield theory	26
2.5.1	Excitation of plasma waves	27
2.5.2	Injection methods	32
2.5.3	Wakefield acceleration	36
2.6	Computational tools	39

3	Kilohertz-rate laser-plasma acceleration	44
3.1	e-KAIO	46
3.1.1	Description and design considerations	46
3.1.2	Experimental testing and verification	50
3.1.3	Summary	57
3.2	Space radiation reproduction	60
3.2.1	Reference model and testing requirements	63
3.2.2	Identification of acceleration regime	64
3.2.3	Characterisation of acceleration regime	70
3.2.4	Stability of the accelerator	76
3.2.5	Electron beam parametric improvement	79
3.2.6	Accelerator flexibility	84
3.2.7	Summary	85
4	Hybrid plasma photocathode acceleration	87
4.1	Simulation parameters and setup	89
4.2	Witness charge is fully decoupled from driver beam charge	94
4.3	Witness energy gain is independent of driver beam energy	97
4.4	Witness energy gain is independent of driver beam energy spread	101
4.5	Beam loading and witness chirp reduction by driver beam depletion	105
4.6	Witness production has high spatio-temporal release tolerance	107
4.7	Witness beam energy gain self-adapts to driver charge density fluctuations	110
4.8	Summary	115
5	Plasma afterglow metrology	121
5.1	Experimental setup	126
5.1.1	Laser system characterisation	126
5.1.2	Beam transport	130

5.1.3	Experimental chamber	134
5.1.4	Laser-gas interaction	137
5.1.5	Afterglow diagnostics	139
5.2	Experimental results - Flooded chamber mode	143
5.2.1	Laser pointing angle variation	145
5.2.2	Laser energy variation	149
5.2.3	Gas density variation	155
5.3	Simulations	159
5.4	Experimental results - Gas jet mode	168
5.5	Summary	173
6	Summary and Outlook	176
	Bibliography	179

List of Figures

2.1	Generalised Gaussian laser beam	10
2.2	Single vs Multi-Photon Ionisation	14
2.3	Tunneling vs Barrier Suppression Ionisation	15
2.4	ADK ionisation rates for various elements	17
2.5	Linear regime of LWFA	30
2.6	Non-linear regime of PWFA	31
2.7	Generalised PIC cycle	40
3.1	3D e-KAIO visualisation	46
3.2	Schematic of the e-KAIO acceleration module	48
3.3	e-KAIO graphical user interface (GUI)	51
3.4	Pumping system gas load measurements	52
3.5	PHASICS gas profile lineout	54
3.6	Coherent Astrella far-field beam profile	56
3.7	Side view shadowgraphy comparison	58
3.8	AE8-MAX trapped electron model	65
3.9	Optimised electron energy spectra - Helium	68
3.10	Optimised electron energy spectra - Argon	69
3.11	Laser pulse parameter evolution along the propagation	71
3.12	Laser pulse E-field spatial and intensity evolution along the prop- agation	73

3.13	Electron momentum after second laser pulse collapse	74
3.14	2D longitudinal wakefield and on-axis longitudinal 1D wakefield prior to the second pulse collapse	75
3.15	Stability of the accelerated electron beam	77
3.16	Accelerated beam profile - Linearly polarized laser	78
3.17	Accelerated beam profile - Circularly polarized laser	80
3.18	Accelerated electron energy spectra - Circularly polarized laser in Argon	81
3.19	Accelerated electron beam divergence comparison	82
3.20	Favourable skin cancer therapy electron energy spectra	84
4.1	Collinear staged plasma photocathode wakefield accelerator in the laboratory frame	89
4.2	PWFA showcase simulation	92
4.3	Witness charge when driver charge or laser energy varies	95
4.4	Impact of driver beam electron energy (0.125–3.5 GeV) on witness energy gain	98
4.5	Impact of driver beam energy spread on witness bunch energy gain	102
4.6	Wakefield evolution showing beam-loading effects from driver en- ergy spread	106
4.7	Intrinsic robustness of witness bunch production against photo- cathode laser spatio-temporal jitter, shown as a function of longi- tudinal and radial release position within the blowout	108
4.8	Intrinsic mitigation of witness energy fluctuation arising from driver density-induced wakefield jitter.	111
5.1	Experimental layout	125
5.2	Laser pulse duration and spectrum	127
5.3	Coherent Legend near-field beam profile	128

5.4	Laser energy stability	129
5.5	Coherent Legend far-field beam profile	131
5.6	Coherent Legend PTFE screen profile	132
5.7	LIDT, B-integral and pulse lengthening clarification	133
5.8	Pumping system performance	135
5.9	Particle-in-cell simulated plasma filament and electron spectra . .	139
5.10	Plasma afterglow - side and top view	140
5.11	Spatial and integrated afterglow spectrum	141
5.12	Afterglow intensity stability	142
5.13	Asymmetric afterglow signal	144
5.14	Spatially mapped ionization regions	145
5.15	Afterglow Images - Laser pointing angle variation	146
5.16	Laser spatial movement scan	147
5.17	Small range laser energy scan	150
5.18	Afterglow Images - Laser energy variation	151
5.19	Laser energy - Afterglow FWHM length correlation	152
5.20	Wide range laser energy scan	154
5.21	Wide range gas density scan	155
5.22	Afterglow Images - Gas density variation	157
5.23	Wide range gas density scan afterglow trends	158
5.24	Linear laser-gas propagation and focusing case	160
5.25	Ionization-induced defocusing laser-gas propagation and focusing case	162
5.26	Simulated laser pulse intensity and waist evolution	163
5.27	Ionization induced defocusing afterglow metrology	165
5.28	Dynamic spatial replenishment propagation case	167
5.29	Afterglow Images - Gas jet mode	168
5.30	Afterglow Images - Gas jet transverse movement variation	170

5.31 Gas nozzle profile reconstruction	172
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List of Tables

3.1	Experimental parameters and their associated acceleration regimes	59
5.1	Camera specifications	126

Chapter 1

Introduction

1.1 Particle accelerators

For over a century, particle accelerators have driven scientific discovery and human progress, enabling researchers to tackle fundamental physics questions such as uncovering and understanding the basic building blocks of matter. Early examples include landmark achievements such as the detection of the electron [1] and the artificial splitting of the atom [2]. Nowadays, accelerators are reaching higher and higher energies, further probing the validity of the Standard Model and beyond; such as the discovery of the Higgs boson [3] and searches for supersymmetry [4]. These high energy particle beams can also emit coherent x-ray radiation in free electron lasers (FEL) [5] and synchrotrons [6], allowing the resolution of ultrafast processes in molecules and atoms on their natural timescales [7]. However, particle accelerators are not restricted to fundamental physics research, and in practice, the majority of accelerator facilities support diverse scientific and technological domains, delivering substantial societal and industrial benefits. Notable applications include radiation therapy for cancer treatment [8], medical imaging [9] and the doping of semiconductors for electronic devices [10].

Highest modernity radio frequency (RF) accelerators such as the Large Hadron

Collider (LHC) & Stanford Linear Accelerator Center (SLAC) employ oscillating electromagnetic fields in a series of RF cavities to incrementally transfer energy to charged particles, thus increasing their velocity [11]. Whilst these accelerators have been successfully refined and deployed over many decades, and their scientific and technological significance is undeniable [12], their performance is ultimately constrained, in large part because they can only sustain maximum accelerating electric fields of ≈ 55 MV/m before dielectric breakdown starts to occur in their cavity walls [13]. This condition dictates that in order to achieve an ever increasing energy budget, RF accelerator systems now require kilometre scale real estate, which is extremely costly to both build and maintain [14]. These restraints significantly limit the number of high-energy particle accelerators, thereby slowing the rate of scientific advancement. It is therefore imperative to explore a number of novel accelerator concepts that can increase the accelerating gradient by orders of magnitude, thus shrinking the size of the accelerator e.g. from the kilometre to the metre scale.

One approach that has seen significant advancements involves adopting the broken-down, fourth state of matter - plasma, as the operating medium of the accelerator. Defined as a quasineutral gas of charged and neutral particles, a plasma doesn't suffer the same limitation as RF cavities since it is already in an ionised state and can support GV/m scale electric fields in regions where charge separation occurs locally. Intense laser pulses or particle beams can excite these plasma structures, referred to as wakes, which effectively serve as microscale accelerating cavities. A charged particle bunch injected at the appropriate phase within the wakefield can undergo rapid energy gain, achieving acceleration over propagation distances that are orders of magnitude shorter than those required in conventional accelerator technologies. Therefore, plasma-based acceleration schemes [15] provide a path towards compact, high-energy particle sources that can be deployed within laboratory-scale infrastructures, as well as a viable approach for creating

next-generation particle colliders and advanced light sources facilities, where the physical scale and financial demands of an equivalent conventional accelerator would render it impractical.

To date, plasma accelerator research has largely focused on the optimisation of key output particle beam parameters such as energy, energy spread, charge, divergence and duration, with developments in laser technology and the discovery of injection schemes that help gain more control over the injection process leading to substantial progress in this area. However, a significant thrust is now being made towards improving the practicality of these setups in order to turn plasma accelerators into actual machines that can be used for a variety of applications. This thesis outlines advancements towards this aim at multiple branches of the plasma accelerator ecosystem, namely utilising higher repetition rate kHz lasers to drive the accelerator, successfully mitigating fluctuations in the output particle beam and establishing a live in-situ non-destructive diagnostic.

1.2 Thesis outline

The presented work is divided into five parts. Chapter 2 introduces the fundamental concepts necessary for understanding the physical processes underlying plasma-based particle acceleration.

Chapter 3 presents developments toward high repetition rate laser–plasma acceleration through two complementary efforts. The first involves proof-of-principle experimental validation of a laser–plasma accelerator module optimised for operation with kHz laser drivers, with a view toward industrial deployment of compact electron-beam sources. The second comprises a simulation-based investigation demonstrating that current kHz laser systems are capable of producing electron fluxes with properties analogous to those encountered in space radiation environments, thereby supporting the growing demand for appropriate

beam sources for radiation-hardness assurance of space-borne electronics.

Chapter 4 lays out a simulation-based study that illustrates how output beam parameters can be effectively stabilised against several inherent experimental fluctuations in a secondary beam-driven plasma photocathode-based wakefield refinement stage.

Chapter 5 details the results of a plasma afterglow metrology experiment, where the relationship between relevant experimental parameters and the afterglow signal was quantified, allowing the light to be exploited as an online non-destructive diagnostic for measuring historically difficult to diagnose quantities at the interaction point. Simulations performed in support of these results are also presented within the chapter.

Chapter 6 concludes the thesis by providing both a high-level synopsis and future outlook for the work conducted throughout this project.

1.3 Role of the author

Results presented in Chapter 3 relating to the e-KAIO laser-plasma accelerator module were achieved alongside researchers from Forschungszentrum Jülich, Heinrich-Heine-Universität, Ecole Polytechnique and SourceLAB. Francois Sylla, Jerome Faure and Markus Büscher conceptualised the module and its associated methodology. The author contributed to the software design, investigation, validation and formal analysis, alongside Christian Greb, Esin Aktan, Viktoria Nefedova, Cedric Sire and Roman Adam. These results have been published in [16].

The author designed the prefile, performed the PIC simulations and analysed the results in the space radiation reproduction study discussed in Chapter 3. These efforts were supported by Arpad Lenart, who has recently conducted an experimental campaign at Heinrich-Heine-Universität that aims at validating

these results.

The numerical simulation work in Chapter 4 was led by David Campbell, who performed the final production FBPIC simulations and analysis to create the figures displayed throughout this chapter. The author contributed to the initial exploratory VSIM simulations and analysis where the underlying physics governing the stabilisation mechanisms was characterised, as well as figure design and conceptualisation in Python. In particular, he led the investigation outlined in Sections 4.3 & 4.4 into how ‘witness’ electron beam energy spread varies w.r.t. driver electron beam energy and energy spread. Thomas Heinemann helped develop the simulation prefile utilised during this study. Thomas Wilson implemented the external laser in FBPIC. These results have been published in [17].

The author carried out the experimental, theoretical and simulation-based research for the work presented in Chapter 5, as well as the data mining and analysis. Grace Manahan aided in the design of the optical beamline and experimental setup. Natascha Thomas helped acquire data when the experiment was placed in gas jet mode. Constantin Aniculaesei performed Ansys Fluent simulations of the gas flow through the nozzle in gas jet mode to facilitate a comparative assessment of our afterglow results.

Chapter 2

Fundamental Concepts

In this chapter, the fundamental concepts necessary to interpret the results presented in this thesis are laid out. Since many of these results deal with interactions between electromagnetic fields and matter, the chapter begins with a brief overview of electrodynamics. Next, the basic concepts of laser physics are introduced, followed by a review of essential plasma physics principles, such as how to generate a plasma, its evolution over time, and its associated recombination features. The specific interaction of an intense laser pulse with plasma is then examined, both in terms of individual electrons and as a collective species, followed by a discussion on the formation of plasma wakes, methods for injecting electrons into them, and the succeeding acceleration process. Finally, the chapter concludes by highlighting computational tools that can be used to model aspects of the physics outlined in this work.

2.1 Electrodynamics

Maxwell's equations [18] describe electromagnetic phenomena, including the generation and propagation of electric ($\mathbf{E}$) and magnetic ($\mathbf{B}$) fields and their interaction with each other, as well as with charges and currents. Their differential

form is shown below

$$\nabla \cdot \mathbf{E} = \frac{\rho}{\epsilon_0} \quad (2.1)$$

$$\nabla \cdot \mathbf{B} = 0 \quad (2.2)$$

$$\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t} \quad (2.3)$$

$$\nabla \times \mathbf{B} = \mu_0 \left(\mathbf{j} + \epsilon_0 \frac{\partial \mathbf{E}}{\partial t} \right) \quad (2.4)$$

In these equations above, the constants μ_0 and ϵ_0 are the magnetic permittivity and electric permittivity in vacuum respectively, which relate to the speed of light c by

$$c = \frac{1}{\sqrt{\epsilon_0 \mu_0}} \quad (2.5)$$

ρ is the charge density

$$\rho = ne \quad (2.6)$$

and $\mathbf{j}$ is the current density

$$\mathbf{j} = en\mathbf{v} \quad (2.7)$$

where e is the elementary charge (1.602×10^{-19} C), n is the charge distribution's number density, and $\mathbf{v}$ is the velocity of the charge distribution.

It is also possible to express the magnetic field using the vector potential ($\mathbf{A}$)

$$\mathbf{B} = \nabla \times \mathbf{A} \quad (2.8)$$

and for the electric field to be expressed using the vector potential and the scalar potential (Φ)

$$\mathbf{E} = -\frac{\partial}{\partial t} \mathbf{A} - \nabla \Phi \quad (2.9)$$

Now, Eq. 2.8 and 2.9 can be substituted into Eq. 2.1-2.4, to give the wave equations for the electric and magnetic fields (in vacuum with no free charges).

$$\left[\nabla^2 - \frac{1}{c^2} \frac{\partial^2}{\partial t^2} \right] \begin{bmatrix} \mathbf{E} \\ \mathbf{B} \end{bmatrix} = \begin{bmatrix} 0 \\ 0 \end{bmatrix} \quad (2.10)$$

One solution to the equations in 2.10 is the plane wave¹, a monochromatic wave of uniform intensity, which takes the form

$$\mathbf{E} = \mathbf{E}_0 \cos [\mathbf{kz} - \omega t + \phi] \quad (2.11)$$

$$\mathbf{B} = \mathbf{B}_0 \cos [\mathbf{kz} - \omega t + \phi] \quad (2.12)$$

where ϕ is the phase of the wave, ω is the angular frequency, $\mathbf{E}_0$ and $\mathbf{B}_0$ giving the amplitude, and $\mathbf{k}$ being the wave vector for the wave travelling in $\mathbf{z}$. The wave vector and the angular frequency are connected via the dispersion relation in vacuum.

$$k = \frac{\omega}{c} \quad (2.13)$$

The force exerted on a point charge q , possessing mass m and velocity $\mathbf{v}$, in the presence of an electromagnetic field is governed by the Lorentz force law [19]

$$\mathbf{F} = \frac{d\mathbf{p}}{dt} = \frac{d\gamma m \mathbf{v}}{dt} = q(\mathbf{E} + \mathbf{v} \times \mathbf{B}) \quad (2.14)$$

where $\gamma = \frac{1}{\sqrt{1-v^2/c^2}}$ is the relativistic factor. Throughout the majority of this thesis, the dynamics of electrons in electromagnetic fields will be analysed, where electrons are characterized by the elementary charge e and mass m_e (9.109×10^{-31} kg).

¹It should be noted that plane waves are purely mathematical and cannot physically exist as they must have a value everywhere in space, which would require infinite energy, however they can be a useful approximation.

2.2 Laser beams

The spatial profile of all laser pulses discussed in this thesis can be characterized by a Gaussian distribution. For a radially symmetric profile with radius r , the electric field of a monochromatic beam near the focal point can be expressed as [20]

$$\mathbf{E}(r, z, t) = \frac{E_0}{2} \underbrace{e^{-\frac{r^2}{w(z)^2}}}_{\textcircled{1}} \underbrace{e^{-\frac{(t-z/c)^2}{\tau_0^2}}}_{\textcircled{2}} \underbrace{\cos(\omega_L t - k_L z + \varphi_L)}_{\textcircled{3}} \mathbf{e}_{pol} + cc., \quad (2.15)$$

where the first term $\textcircled{1}$ describes the radial envelope with the transverse $1/e$ radius of the beam $w(z)$, the second part $\textcircled{2}$ represents the temporal envelope with τ_0 as the pulse duration defined at the $1/e$ height of the electric field (see Figure 2.1 (b)), and $\textcircled{3}$ is a rapidly oscillating phase term containing the carrier angular frequency ω_L . Finally, $\mathbf{e}_{pol}$ is the vector describing the polarization of the laser pulse in the x-y plane i.e. $\mathbf{e}_{pol} = \mathbf{e}_{x,y}$ for linear and $\mathbf{e}_{pol} = (\mathbf{e}_x \pm i\mathbf{e}_y)/\sqrt{2}$ for circular polarization.

The beam's transverse radius changes along the propagation direction according to

$$w(z) = w_0 \sqrt{1 + \left(\frac{z}{z_R}\right)^2} \quad (2.16)$$

where w_0 is the radius at the $1/e$ height of the maximum electric field and z_R is the Rayleigh length

$$z_R = \frac{\pi w_0^2}{\lambda_L} \quad (2.17)$$

which is defined as the distance from the beam waist to the position along the z-axis where the intensity drops to half its maximum value, with λ_L denoting the laser wavelength.

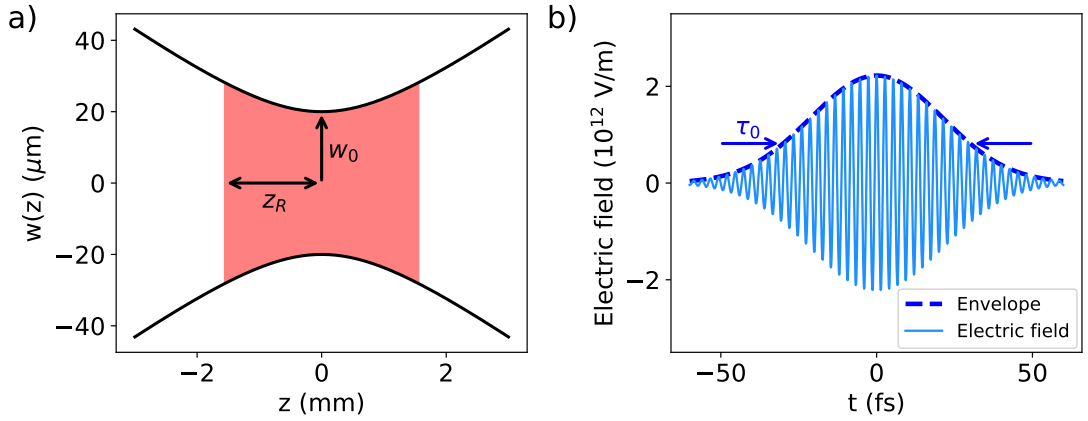


Figure 2.1: Generalised Gaussian laser beam. a) The beam waist $w(z)$ evolution along the propagation z as indicated by the black lines, and the interaction region ($-z_R \leq z \leq z_R$) visualised by the red shaded area. b) The temporal profile of the focused pulse (light blue) and its associated envelope (dark blue). The arrows mark the $1/e$ height pulse duration of the electric field.

For a typical Ti:sapphire laser with a wavelength of 800 nm, focusing to $w_0 = 20 \mu\text{m}$ yields a Rayleigh range of approximately 1.6 mm. This scenario is illustrated in Figure 2.1 (a), where w_0 and z_R are represented by the two black arrows. Utilising the envelope of a Gaussian laser pulse, as shown in 2.1 b), is a common approximation technique that helps to simplify calculations such as cycle-averaged tunneling rates for ionization (see Section 2.3.1), or reduce the computational costs associated with running PIC simulations (see Section 2.6).

To determine how close a real laser beam is to an ideal Gaussian beam, the dimensionless M^2 or beam quality parameter was introduced [21] to quantify the beam's deviation from the theoretical minimum beam divergence and waist size. It is defined as

$$M^2 = \frac{\pi w_0 \theta}{\lambda_L} \quad (2.18)$$

where θ is the beam divergence angle. A perfect Gaussian beam has $M^2 = 1$ meaning real laser beams have $M^2 > 1$, with lower M^2 values signifying better

beam quality. The M^2 parameter limits the degree to which a beam can be focused as well as impacting how it propagates, making it a key factor for laser plasma accelerators. The laser's beam waist w_0 and pulse duration τ_0 are related to the measured FWHM values of the spatial and temporal intensity profiles through the spot size, $d_{FWHM} = \sqrt{2\ln(2)}w_0$, and pulse duration, $t_{FWHM} = \sqrt{2\ln(2)}\tau_0$. Careful attention must be paid to which expression is required when attempting to model effects such as field ionization (see Section 2.3.1), as although both definitions are valid, inconsistent use can result in markedly different values from Equation 2.15.

From the d_{FWHM} , t_{FWHM} and the measured total laser pulse energy value W_L , the laser peak intensity can be determined

$$I_0 = 0.83 \frac{W_L}{t_{FWHM} d_{FWHM}^2} \quad (2.19)$$

which in turn allows the peak electric field E_0 to be expressed in terms of I_0

$$E_0 = \sqrt{\frac{2I_0}{c\epsilon_0}} \quad (2.20)$$

Modern high-intensity laser systems routinely achieve intensities exceeding 10^{19} Wcm^{-2} , corresponding to electric field strengths on the order of 10^{12} Vm^{-1} . These fields surpass atomic binding fields by several orders of magnitude, and therefore result in ionisation and plasma formation within the interaction region.

In the context of this thesis, it is often advantageous to characterize the laser by a unit-less strength parameter, the peak normalized vector potential a_0 [22], which serves as a measure of the relativistic intensity regime relevant to laser-matter interactions. It can be calculated using the following expression

$$a_0 = \frac{e|\mathbf{E}_0|}{m_e\omega_L c} = 0.85 \sqrt{I_0 [10^{18} \text{ W cm}^{-2}]} \lambda_L [\mu\text{m}] \quad (2.21)$$

When $a_0 > 1$, the laser-matter interaction enters a highly relativistic regime,

exemplified by phenomena such as the bubble regime in laser-driven wakefield acceleration [23]. In contrast, $a_0 < 1$ corresponds to a non-relativistic, classical interaction, where, for example, the momentum imparted to electrons in the tunneling ionization regime (see Section 2.3.1) is minimal, facilitating the formation of cold, spatially localized plasma.

2.3 Plasma physics

Plasmas, often described as the fourth state of matter, are composed of an ionized gas of electrons and ions. A common way to characterize plasmas is by examining their response to perturbations [24]. For instance, displacing electrons slightly from the stationary ion background creates a region of charge separation, generating an electric field that depends on both the ion and electron densities, n_i and n_e respectively, which are equal under quasi-neutral conditions, $n_i = n_e$. As the electrons move back toward their equilibrium positions, they undergo oscillations at the plasma frequency

$$\omega_p = \sqrt{\frac{n_e e^2}{m_e \epsilon_0}} \quad (2.22)$$

for which there is an associated plasma wavenumber

$$k_p = \frac{\omega_p}{c} = \sqrt{\frac{n_e e^2}{m_e \epsilon_0 c^2}} \quad (2.23)$$

The inverse of this k_p^{-1} defines the distance light can penetrate into the plasma before its amplitude is significantly attenuated, known as the skin depth, which is fundamentally linked to the plasma wavelength

$$\lambda_p = \frac{2\pi c}{\omega_p} \quad (2.24)$$

the length over which electrons in the plasma oscillate in response to a disturbance.

2.3.1 Plasma generation

Optical ionisation

The creation of a plasma requires at least the partial ionization of atoms or molecules to generate an ensemble of charged particles. Optical ionisation is the primary plasma generation method utilised in this work and it involves electromagnetic radiation exciting one or more electrons from their bound states, as defined by their ionization energy E_{ion} , into free, unbound states. In the field of plasma acceleration, the radiation source is typically a high-intensity laser system (see Section 2.2), capable of ionizing large volumes of gas or solid material.

The laser field possesses a pondermotive potential $\Phi_{pond} = e^2 E_0^2 / 4m\omega_L^2$ and supplies photons of energy $E = \hbar\omega_L$ and frequency ω_L . Here, Φ_{pond} is an expression of the effective quiver energy acquired by an electron oscillating at ω_L in the laser's electromagnetic field, E_0 is the peak electric field amplitude (see Equation 2.20) and $\hbar$ is the reduced Planck's constant.

Ionisation can be induced either by the electromagnetic field itself, or by individual or multiple photons, with the Keldysh adiabaticity parameter [25]

$$\gamma \sim \sqrt{\frac{E_{ion}}{\Phi_{pond}}} \quad (2.25)$$

distinguishing between these various regimes. As a rule of thumb, $\gamma \gg 1$ corresponds to the well-known process of photoionisation in single atoms by one or multiple photons. When the photon energy E exceeds the ionisation energy E_{ion} , a single photon is sufficient to ionise the atom. While this scenario is conceptually straightforward, it poses significant challenges in practical applications, especially when dealing with tightly bound electrons. For instance, hydrogen and helium, commonly used in plasma accelerators, require ionisation energies of 13.6 eV (corresponding to ~ 91 nm photons for atomic hydrogen) and 24.6 eV (~ 50 nm photons for helium), placing them in the extreme ultraviolet (EUV) range.

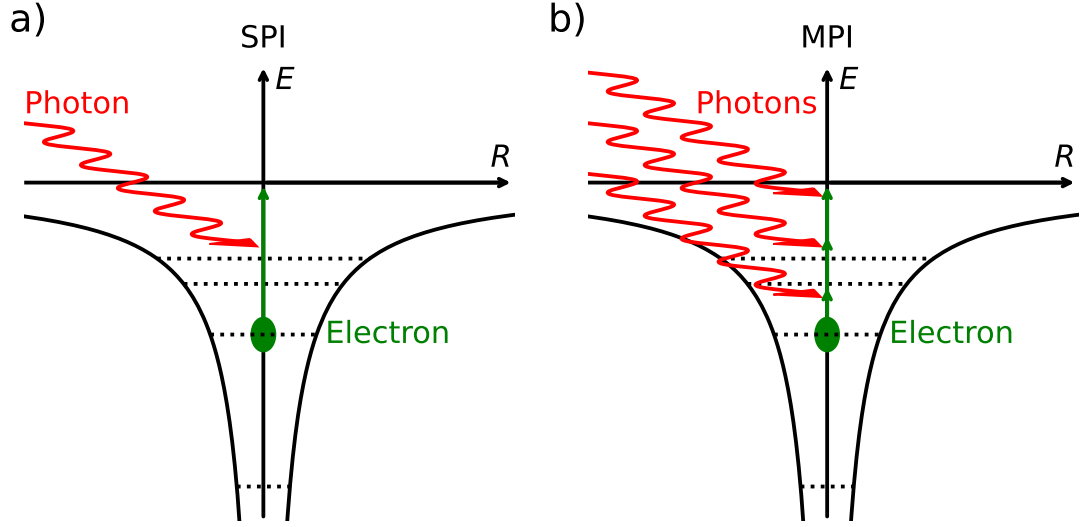


Figure 2.2: Single vs Multi-Photon Ionisation. a) Single-Photon Ionization - One photon delivers enough energy to free the electron from its bound state. b) Multi-Photon Ionization - Several low-energy photons arrive simultaneously at the atom, and together their combined energy is sufficient to release the electron.

However, powerful laser sources at these wavelengths are not readily available, and instead radiation sources must rely on inefficient techniques such as high harmonic generation [26].

Alternatively when $E < E_{ion}$, several lower frequency photons can simultaneously interact with a single atom, collectively transferring their energy to induce ionisation. The likelihood of this scheme [27] depends strongly on the intensity I_L and frequency of the incoming light, with the ionisation probability P scaling strongly with the number of photons n involved $P \propto I_L^n$. TW-class laser systems can comfortably achieve the intensities ($> 10^{10} \text{ Wcm}^{-2}$) required to efficiently drive this process. Both single and multi-photon ionisation mechanisms are illustrated in Figure 2.2.

However, most modern high-intensity lasers operate in a regime where the ponderomotive potential of the laser field far surpasses the atomic binding potential $\gamma \ll 1$, resulting in the Coulomb barrier being suppressed below the binding

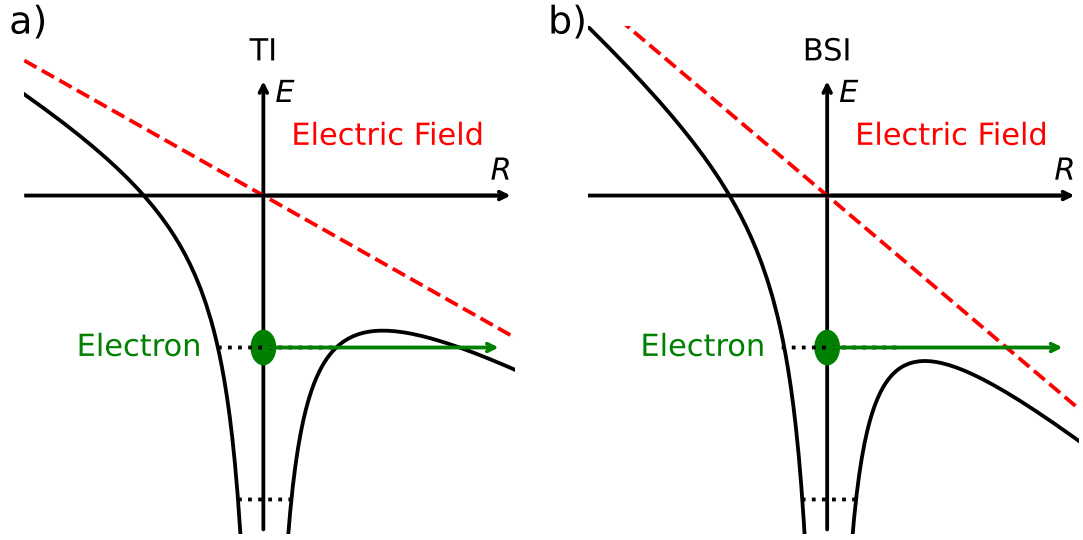


Figure 2.3: Tunneling vs Barrier Suppression Ionisation. a) Tunneling Ionisation - A strong external electric field lowers the effective potential barrier of the atom, thereby enhancing the quantum mechanical tunneling probability for the electron to escape. b) Barrier Suppression Ionisation - An intense external electric field significantly distorts the atomic binding potential, suppressing it over a sufficient duration to allow the initially bound electron to escape classically.

energy and the electron escaping through classical motion (see Figure 2.3 b)). This phenomenon, known as Barrier Suppression Ionisation (BSI) [28], occurs when the laser's electric field strength reaches a critical threshold

$$\epsilon_c = \frac{E_{ion}^2}{4Ze^3} \quad (\text{Gaussian units}) \quad (2.26)$$

where Z is the atomic number and e is the elementary charge. Equating this to the laser's peak electric field E_0 yields an effective appearance intensity

$$I_{app} = \frac{c}{8\pi} \epsilon_c^2 = \frac{cE_{ion}^4}{128\pi Z^2 e^6} \quad (\text{Gaussian units}) \quad (2.27)$$

which when simplified can be expressed as

$$I_{app} \simeq 4 \times 10^9 \left(\frac{E_{ion}}{\text{eV}} \right)^4 Z^{-2} \text{ Wcm}^{-2} \quad (2.28)$$

Looking at Hydrogen which is the simplest example, where $Z = 1$ and $E_{ion} = 13.61$ eV, $I_{app} \simeq 1.4 \times 10^{14}$ Wcm $^{-2}$. When $I > I_{app}$, liberated electrons promptly interact with the laser field, resulting in rapid acceleration and the generation of high-energy (hot) plasma populations. At even greater intensities, the interaction can excite non-linear plasma waves, manifesting as propagating density modulations. This regime is characteristic of laser wakefield acceleration (LWFA), wherein the laser pulse efficiently ionizes the gaseous target, often to full depletion of bound electrons, and then accelerates a portion of the free electrons. While beneficial for LWFA and related ultrafast acceleration schemes, such conditions are unfavourable for applications requiring controlled, level-specific ionization or the preservation of cold, quasi-static plasma environments.

Finally, there is another regime that exists in which $\gamma < 1$ but the field amplitude is strong enough to significantly alter the potential. In this instance, the electric field distorts the Coulomb potential such that an electron finds itself close to the finite potential barrier. The quantum nature of the electron then allows for tunnelling through this reduced barrier to take place (see Figure 2.3 a)), in a process termed Tunneling Ionisation (TI). The tunneling ionization rate can be expressed semi-classically by averaging the electric field over a single optical cycle, as laid out in the widely adopted Ammosov–Delone–Krainov (ADK) theory [29]. This approach assumes that the tunneling time, governed by the electron’s bound-state dynamics, is much shorter than the laser field’s oscillation period, thereby justifying an adiabatic approximation [30]. In practical terms, the ADK formalism yields a rate expression that effectively estimates the ionization probability of valence electrons under strong-field conditions [31]

$$W_{ADK} \approx 1.52 \times 10^{15} \frac{4^{n^*} E_{ion}(\text{eV})}{n^* \Gamma(2n^*)} \left(20.5 \frac{E_{ion}^{3/2}(\text{eV})}{E_0(\text{GVm}^{-1})} \right)^{2n^*-1} \times \exp \left(- 6.83 \frac{E_{ion}^{3/2}(\text{eV})}{E_0(\text{GVm}^{-1})} \right) \quad (2.29)$$

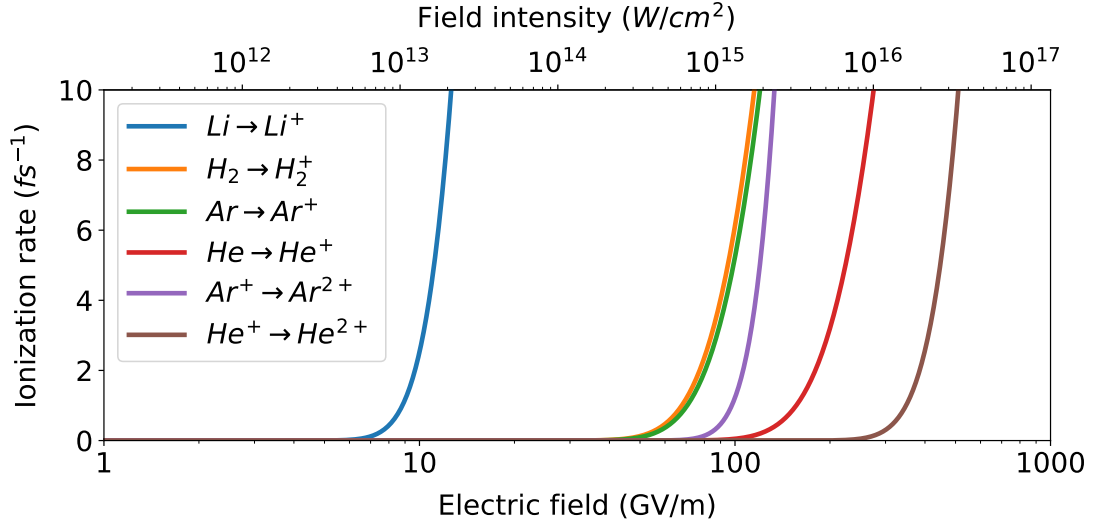


Figure 2.4: ADK ionisation rates for various elements. The ionisation rates, as calculated by Equation 2.29, for a selection of ionisation levels of gases most relevant to the ensuing work in this thesis, shown as a function of both the electric field and field intensity.

where Γ is the gamma function and $n^* \approx 3.69Z E_{ion}^{-1/2}(\text{eV})$ is the effective principal quantum number of the atom or ion's ground state. This equation is valid as long as the laser's electric field does not exceed $E_0 = (\sqrt{2} - 1)E_a(E_{ion}(\text{eV})/27.2\text{eV})^{3/2}$ where $E_a = (4\pi\epsilon_0)^{-3}m_e^2e^5/\hbar^4 \approx 5.14 \times 10^{11} \text{ Vm}^{-1}$. At amplitudes exceeding this critical value, BSI becomes the dominant process over TI. Tunneling ionization predominantly liberates valence electrons and, when driven by intensities just sufficient to ensure consistent ionization rates, enables the controlled generation of cold plasma distributions. In gases or gas mixtures with well-separated ionization potentials, this mechanism can selectively ionize a specific valence electron population while leaving other electronic states largely unaffected, as demonstrated by Figure 2.4. Such selective and minimally perturbative ionization makes tunneling ionization particularly well-suited for applications requiring high fidelity.

Impact ionisation

Outside of optically generating a plasma, collisions between energetic particles

(e.g. electrons or ions) and neutral atoms or molecules, constitutes another fundamental mechanism for plasma creation, particularly in high-density or thermally driven environments. Impact ionization requires that incident particles possess kinetic energy exceeding the ionization threshold, i.e. $W_{kin} > E_{ion}$. The resultant increase in the particle density n_p can be quantitatively described by a rate equation governing binary collisional processes [32]

$$\frac{\partial n_p}{\partial t} = n_i n_g \sigma_i v \quad (2.30)$$

where a distribution of incident electrons or ions with density n_i , propagating with velocity v , traverses a gaseous medium characterized by a target atom or molecule density n_g . The interaction probability is governed by the ionization cross section σ_i , which is a function of the velocity v of the incoming particles and represents the statistical likelihood of ionizing scattering events transpiring.

It is noteworthy that the velocity dependence in Equation 2.30 encapsulates two counteracting phenomena: while higher-velocity particles traverse the target medium more rapidly, thereby increasing the frequency of encounters with atomic or molecular constituents, their corresponding ionization cross section is significantly reduced compared to that of electrons with kinetic energies just above the ionization potential. Impact ionization is an essential procedure for plasma generation in electric discharge devices, which operate by sustaining current through a gaseous media [33]. In the low-current Townsend regime, externally applied electric fields accelerate arbitrarily emitted electrons and ions, which subsequently undergo inelastic collisions with neutral gas atoms or molecules. These interactions produce additional charge carriers that are likewise accelerated by the field, initiating an electron avalanche. This self-sustaining ionization process leads to continual plasma formation, which can be harnessed for various applications, most notably the excitation of characteristic optical emissions in gas-discharge lamps.

2.3.2 Ionisation-induced defocusing

As established in Section 2.3.1, the majority of gases will undergo multiple ionization at laser intensities exceeding 10^{16} Wcm^{-2} . This might suggest that plasma formation should occur within the leading edge of a TW laser pulse, thereby allowing the bulk of the pulse to propagate through a fully ionized, homogeneous plasma column. However, this idealized picture is disrupted by the phenomenon of ionization-induced defocusing [34, 35]. Near the rising edge of the pulse, where the electric field amplitude approaches the ionization threshold, the gas along the central beam axis ($r=0$) will be ionised more, creating a steep radial electron density gradient. Additionally, this varying ionization rate also applies in the temporal domain (along z), resulting in spectral blue-shifting of portions of the transmitted pulse [36]. The refractive index of the plasma generated via ionization is given by

$$\eta(r, t) = \left(1 - \frac{n_e(r, t)}{n_c}\right)^{1/2} \quad (2.31)$$

where $n_e(r, t)$ is the local electron density and n_c is the critical density (see Section 2.4). An increased electron density at the beam centre results in a local minimum of the refractive index, effectively creating a defocusing plasma lens for the trailing portion of the pulse. Consequently, at high gas pressures, the laser beam undergoes significant deflection prior to reaching its geometrical focus resulting in drastically reduced peak intensities [37]. For a laser spot size d_{FWHM} , the beam deflection initiated by ionization-induced defocusing scales as

$$\theta_I \sim \frac{1}{d_{FWHM}} \int \frac{n_e(0)}{n_c} dz \quad (2.32)$$

quantifying that light rays are refracted away from regions of elevated electron density. Conversely, a Gaussian beam focused in vacuum is limited solely by diffraction and adheres to the following relation [38]

$$\theta_D = \frac{d_{FWHM}}{z_R} \quad (2.33)$$

Comparing θ_D and θ_I across identical propagation lengths, we can define when ionization-induced defocusing becomes the dominant deflection effect

$$\frac{n_e}{n_c} > \frac{\lambda}{\pi z_R} \quad (2.34)$$

If ionization is not yet complete at this stage, the plasma density reaches an effective saturation point, as further focusing of the beam can no longer occur. To mitigate the effects of ionization-induced defocusing, high-intensity laser experiments are typically conducted either in pre-ionized plasmas [39] or using gas jet targets, where the laser is focused in vacuum prior to interacting with the neutral gas medium.

2.3.3 Plasma evolution

Following its generation, the plasma evolves on nanosecond timescales, which are considerably longer than the transit time of a laser or particle beam through a laboratory-scale plasma. Nonetheless, characterising this temporal evolution is critical for deciphering the plasma state that will be encountered by a probe or drive beam that arrives with a significant delay relative to the initial plasma formation. The dynamics of this evolution are governed primarily by the initial electron temperature and plasma density, but can be further influenced by post-ionisation energy deposition mechanisms. These include a range of coupled, non-linear processes including further collisional ionisation [40, 41], bremsstrahlung emission [42] and hydrodynamic expansion [43], where the latter may involve the emergence of shock fronts, which can propagate through and interact with the ambient medium, driving secondary ionisation and modifying the local plasma conditions. Accurate modelling of such behaviour typically necessitates the use

of magneto-hydrodynamic (MHD) simulations, which are highly sensitive to the inputted original plasma parameters and external boundary conditions.

Recombination of electrons and ions to form neutral atoms continuously competes with ionisation processes, and in the absence of sustained external energy, recombination will ultimately prevail, leading to the complete neutralisation of the plasma. Electrons and ions may follow multiple routes towards a fully recombined state, with the most prominent mechanism being radiative recombination



where an ion A^+ captures an electron e^- to form a neutral atom A , with the excess kinetic energy greater than the atomic binding energy emitted as a photon $\hbar\omega$. Alternatively, the excess energy can be transferred to a second free electron in an effect known as three-body recombination



The dominant recombination pathway is governed by the reaction cross-section, which generally depends on the ion's electronic structure and the plasma temperature. It is important to note that recombination rates are typically enhanced at lower kinetic energies of the reacting particles [44]. Neutral atoms formed via recombination processes may retain bound electrons in excited states, which will subsequently decay to lower energy levels through the emission of photons and cause the plasma to radiate light with distinctive spectral lines of the chosen gaseous medium [45]. This coupled with light emitted from various other excitation/relaxation processes and potentially bremsstrahlung radiation, depending on the temperature plasma electrons are spawned with, constitutes the plasma 'afterglow' [46–48]. The signal is highly sensitive with multi-parameter dependency and can therefore be exploited as a versatile diagnostic able to gauge numerous parameters of the interaction (see Section 5).

2.4 Laser propagation in a plasma

State-of-the-art laser systems, with their high peak intensities, enable the leading edge of the laser pulse to ionize the target medium, allowing the peak of the pulse to interact with a preformed plasma. The resulting ponderomotive force can expel electrons from their equilibrium positions, generating macroscopic regions dominated by space-charge effects. To accurately describe these dynamics, it is essential to analyse the propagation of the laser's electromagnetic wave within a plasma medium.

In Equation 2.13, the dispersion relation in vacuum was given. The dispersion of an electromagnetic wave with frequency ω_L , and wave vector, k_L , in a plasma

$$\omega_L^2 = \omega_p^2 + c^2 k_L^2 \quad (2.37)$$

is altered by the interaction with plasma particles. For $\omega_L = \omega_p$, where ω_p is the plasma frequency, a critical density can be defined

$$n_{e,c} = \frac{\epsilon_0 m_e}{e^2} \omega_L^2 = \frac{1.1}{\lambda_L^2 [\mu\text{m}]} 10^{21} \text{cm}^{-3} \quad (2.38)$$

In plasmas with electron densities $n_e > n_{e,c}$, the plasma frequency exceeds the frequency of the incident electromagnetic wave, enabling electrons to respond effectively to the field oscillations. This leads to strong shielding of the wave within the medium and results in reflection; such plasmas are termed overdense. Conversely, when $n_e < n_{e,c}$, the electron response is insufficient to screen the incoming radiation, permitting the electromagnetic wave to propagate through the medium. In this case, the plasma is referred to as underdense. For a Ti:Sa laser, the critical density is $1.74 \times 10^{21} \text{ cm}^{-3}$. The corresponding group and phase velocities, v_g and v_{ph} , of the electromagnetic wave are given by the refractive index η and Equation 2.37:

$$v_g \equiv \frac{\partial \omega_L}{\partial k_L} = \eta c \quad v_{ph} \equiv \frac{\omega_L}{k_L} = \frac{c}{\eta} \quad \text{where } \eta = \sqrt{1 - \frac{\omega_p^2}{\omega_L^2}} < 1 \quad (2.39)$$

The phase velocity of the laser is determined by its driver - specifically, by the group velocity of the pulse. This phase velocity always exceeds the speed of light c , whereas the group velocity, which represents the laser's actual propagation speed through the plasma, is always less than c . The relativistic Lorentz factor associated with the light wave is defined as follows:

$$\gamma_p = \frac{1}{\sqrt{1 - (\frac{v_g}{c})^2}} \approx \frac{\omega_L}{\omega_p} \quad (2.40)$$

2.4.1 Relativistic self-focusing

In the relativistic regime ($a_0 > 1$), the laser-plasma interaction induces modifications to both the local electron density and the effective electron mass via the non-linear ponderomotive force. These perturbations alter the local refractive index, thereby introducing a phase modulation across the pulse envelope and causing intensity-dependent propagation characteristics. As the trajectory of the electrons in the presence of a strong laser field is mainly determined by their quiver motion $\gamma = \sqrt{1 + a^2}$, where a is the local, instantaneous normalized vector potential, the revised refractive index can be approximated as [49]

$$\eta(r, z) \approx 1 - \frac{1}{2} \left(\frac{\omega_p}{\omega_L} \right)^2 \left(1 - \frac{a^2(r)}{2} + \frac{\delta n_e(r)}{n_{e,0}} + \frac{\Delta n_{ext}(r)}{n_{e,0}} \right) \quad (2.41)$$

In Equation 2.41, the second term $\frac{\delta n_e(r)}{n_{e,0}}$ accounts for the radial modulation of the plasma density induced by the ponderomotive force, which drives transverse electron expulsion and consequently facilitates ponderomotive self-focusing. The third term $\frac{\Delta n_{ext}(r)}{n_{e,0}}$ represents a pre-formed plasma channel, typically established via an external mechanism such as an axial electrical discharge or a preceding

auxiliary laser pulse, and provides an externally imposed refractive index profile conducive to guiding. The first term $\frac{a^2(r)}{2}$ encapsulates the influence of the relativistic correction to the electron mass via γ_p . For an on-axis intensity maximum $\frac{\partial a^2}{\partial^2 r} < 0$, this leads to a localized decrease in the plasma refractive index $\frac{\partial \eta}{\partial r} < 0$, resulting in relativistic self-focusing (RSF). This relativistic modification counterbalances natural diffraction, enabling the laser pulse to maintain transverse confinement and propagate with sustained high intensity over distances significantly exceeding the Rayleigh length. The onset of RSF is governed by a critical balance between diffraction and nonlinear self-focusing, leading to the definition of a critical power threshold for stable self-guiding [50, 51]

$$P_c = 17.4 \frac{\omega_L^2}{\omega_p^2} [\text{GW}] \quad (2.42)$$

When the conditions for SF are satisfied, Equation 2.16 which details the beam's transverse evolution along the propagation must be adjusted to

$$w(z) = w_0 \sqrt{1 + \left(1 - \frac{P}{P_c}\right) \frac{z^2}{z_R^2}} \quad (2.43)$$

In a plasma, a laser pulse with power P and waist w will undergo linear diffraction when $P < P_c$, achieve self-consistent transverse confinement for $P = P_c$ and experience non-linear self-focusing for $P > P_c$. In the case of a Ti:Sa laser in a $n_p = 10^{17} \text{ cm}^{-3}$ dense plasma, the critical power is $P_c \approx 303.2 \text{ TW}$. The onset of uncontrolled or 'catastrophic' self-focusing is mitigated by higher-order non-linear effects [52], which act to stabilize the propagation. Experimental observations have confirmed the occurrence of relativistic self-guiding over propagation distances spanning several Rayleigh lengths [53]. Optimal guiding is achieved when the laser spot size is matched to the characteristic blowout radius $w_0 \simeq r_b$, thereby maximizing the overlap between the laser field and the plasma response.

2.4.2 Self-phase modulation

The longitudinal variation of the refractive index can be expressed analogously to its transverse dependence (see Equation 2.41). Utilizing the co-moving coordinate $\xi = z - ct$, the refractive index along z takes the form [50]

$$\eta(\xi) \approx 1 - \frac{1}{2} \left(\frac{\omega_p}{\omega_L} \right)^2 \left(1 - \frac{a^2(\xi)}{2} + \frac{\delta n_e(\xi)}{n_{e,0}} - 2 \frac{\delta \omega(\xi)}{\omega_L} \right) \quad (2.44)$$

Here, the first term $\frac{a^2(\xi)}{2}$ corresponds to the longitudinal gradient of the laser intensity, while the final term $\frac{\delta \omega(\xi)}{\omega_L}$ captures the variation of the laser frequency across the pulse envelope, characterizing its spectral evolution or ‘chirp’. The frequency modulation is governed by the local variation of the refractive index according to [50]

$$\frac{1}{\omega_L} \frac{\partial \omega}{\partial \tau} = - \frac{1}{\eta^2} \frac{\partial \eta}{\partial \xi} \quad (2.45)$$

In the case of Gaussian laser pulses, the intensity gradients at the leading and trailing edges induce red and blue frequency shifts respectively, resulting in spectral broadening - a phenomenon referred to as self-phase modulation (SPM). The associated variation in plasma density leads to a spatially dependent refractive index. In typical laser wakefield acceleration (LWFA) configurations, where the laser pulse encounters a positive plasma density gradient, the resulting group velocity dispersion leads to differential propagation: the pulse front experiences a reduction in group velocity, while the rear is accelerated. This asymmetric group velocity distribution induces temporal compression of the pulse, a behaviour that has been experimentally observed [54, 55] and can be modelled as

$$\tau_{comp} = \tau_0 - \frac{n_{e,0} l}{2c n_{e,c}} \quad (2.46)$$

2.5 Plasma wakefield theory

As early as the 1950s, it was recognized that the strong electric fields sustained in a plasma could, in principle, be harnessed to accelerate charged particles [56, 57]. In 1979, Tajima and Dawson conceptualised a novel scheme in which an intense laser pulse could drive plasma oscillations via the ponderomotive force, generating a trailing wakefield capable of accelerating electrons to relativistic energies over very short distances [58]. Due to the initial difficulty in producing laser pulses of sufficiently short duration to excite such wakefields, a variety of alternative approaches were explored to achieve the necessary laser intensities [59, 60]. The eventual development of compact, high-intensity, short-pulse laser systems [61] made it possible to experimentally realize laser wakefield acceleration (LWFA), enabling the generation and acceleration of high-quality electron beams, commonly referred to as witness or trailing bunches, within the plasma wave [62–64].

In parallel, an alternative approach to driving plasma wakefields was proposed, utilizing high-intensity charged particle beams instead of laser pulses [65]. This beam-driven plasma wakefield acceleration (PWFA) scheme was first experimentally realized in 1988 [66] and has since seen significant advancements, including key milestones such as the energy doubling of a 42 GeV electron beam [67]. Most PWFA implementations employ an electron bunch as the driver to excite a wakefield that subsequently accelerates a trailing electron bunch. However, successful demonstrations have also been achieved using proton beams as drivers for electron acceleration [68], as well as positron drivers for accelerating positron bunches [69].

To date, the development of plasma wakefield accelerators (PWFAs) has been largely limited by the scarcity of high peak-current drive beams, restricting experimental efforts to a small number of specialized radio-frequency (RF) accelerator

facilities. However, recent advancements in laser wakefield acceleration (LWFA) have enabled the generation of high-energy [70], high-charge (approaching the nanocoulomb level) [71, 72], short-duration electron beams [73, 74] with relatively broad energy spreads. These beams, now routinely produced at numerous high-power laser facilities worldwide, possess characteristics that make them well-suited as drivers for PWFAs [75]. Consequently, hybrid LWFA–PWFA (LPWFA) schemes are actively being explored to leverage the complementary advantages of both methods within a compact setup [76–80]. These hybrid systems not only offer a pathway to high-brightness electron beam sources for applications requiring exceptional beam quality but also provide an accessible platform for studying PWFA physics and validating theoretical models at a pace and scale previously unattainable.

The following sections review the fundamental theory underlying plasma wakefield acceleration, beginning with the excitation of plasma waves and proceeding to the injection and acceleration of particle bunches within these structures. While much of the core theoretical framework is applicable to both laser wakefield acceleration (LWFA) and beam-driven plasma wakefield acceleration (PWFA), the emphasis here is placed on the laser-driven case, with comparisons to PWFA provided where relevant.

2.5.1 Excitation of plasma waves

Plasma waves propagate with a phase velocity v_p equal to that of their driver, v_d . In the case of LWFA, the laser pulse travels at the group velocity in plasma (see Section 2.4), which results in $v_{ph} = v_d = c\sqrt{1 - \frac{\omega_p^2}{\omega_L^2}} < c$. This stands in contrast to beam-driven PWFA, where the driver is typically ultra-relativistic, leading to $v_{ph} = v_d \approx c$ in a plasma of uniform density. The phase velocity of the wake has important consequences for both particle injection and acceleration, which are discussed in detail in Sections 2.5.3 and 2.5.2, respectively. The length of

the wake is governed by n_e and λ_p . Ideally, the length of the driver beam σ_z or the laser pulse duration τ_0 should equate to half of the plasma wavelength i.e. $c\tau_0 \approx \lambda_p/2$. This condition ensures optimal excitation of the plasma wave, akin to resonant driving in a harmonic oscillator like a pendulum, where the driver excites the wave up to the zero-crossing point of its oscillation. Expressed in terms of the plasma wavenumber k_p , this ideal situation is typically written as $k_p\sigma_z < 1$.

An excited plasma wave can be described by the unperturbed plasma electron density n_0 , the longitudinal electric field E_z , or the normalised scalar potential $\phi = e\Phi/m_e c^2$. These values are related through $\mathbf{E} = -\nabla\Phi$ and Gauss' law $\nabla \cdot \mathbf{E} = -4\pi e(n - n_0)$, where n is the redistributed plasma electron density post-interaction. The wave equation for the density perturbation caused by the laser's ponderomotive force is [81]

$$\left(\frac{\partial^2}{\partial t^2} + \omega_p^2\right)\frac{\delta n}{n_0} = c^2 \nabla^2 \frac{a^2}{2} \quad (2.47)$$

where the density perturbation $\delta n = n - n_0$ and $\mathbf{F}_p \propto c^2 \nabla(a^2/2)$ is the ponderomotive force. A PWFA beam-driven wake satisfies the same equation for the plasma oscillation, however $\nabla^2 a^2$ is substituted out for the beam charge density driver term n_b . A solution to the wave equation in the linear regime ($a < 1$) is [82]

$$\delta n/n_0 = (c^2/\omega_p) \int_0^t dt' \sin[\omega_p(t - t')] \nabla^2 a^2(\mathbf{r}, t')/2 \quad (2.48)$$

The sinusoidal form reinforces the optimal length condition such that wake-field excitation is most efficient when the laser envelope length characterizing the extent of a^2 is on the order of λ_p . In order to understand the wake excitation in the non-linear regime ($a > 1$), a 1D limit (a large transverse laser spot size) is considered. Assuming the laser pulse is slowly evolving and the phase velocity of the wave is near the speed of light ($\gamma_p^2 \gg 1$), the wake potential can be described

by [49]

$$k_p^{-2} \frac{\partial^2 \phi}{\partial \xi^2} = \frac{(1 + a^2)}{2(1 + \phi^2)} - \frac{1}{2} \quad (2.49)$$

This equation is generally valid for all values of a . As in the linear case, an equivalent driver term for a non-linear PWFA wave can be approximated $n_b \sim \frac{1}{2} \frac{d^2 a^2(\xi)}{d\xi^2}$. It is important to note that in PWFA, the plasma wave is excited more efficiently by the unipolar coulomb force of the charged particle beam instead of the oscillating pondermotive force of an intense laser pulse in LWFA. Advantageously, this allows PWFA to reach the non-linear “blowout” regime at lower electric fields due to its field vector pointing outwards in all directions, enabling greater flexibility when selecting the plasma medium and injection scheme.

In the linear regime, the plasma wave is sinusoidal and the plasma wavelength is expressed by Equation 2.24. In the non-linear regime, the density profile steepens and the electric field becomes a sawtooth shape. The plasma wavelength in the non-linear regime λ_{NP} is larger than that in the linear regime and depends on the wake amplitude

$$\lambda_{NP} = \lambda_p \begin{cases} 1 & \text{for } E_{max}/E_{wb} \ll 1 \\ (2/\pi) E_{max}/E_{wb} & \text{for } E_{max}/E_{wb} \gg 1 \end{cases} \quad (2.50)$$

where E_{max} , equal to $E_{max}(n_e, a_0) = E_{wb}(n_e) \frac{a_0^2/2}{\sqrt{1+a_0^2/2}}$ for the laser-driven case, is the maximum electric field that can be excited in the 1-D limit for a given plasma density and driver strength, and E_{wb} is the cold, non-relativistic wave breaking field (see Equation 2.51).

To better demonstrate 3D wake excitation in the linear and non-linear regimes, simulations using particle-in-cell (PIC) code FBPIC [83] (see Section 2.6) were carried out. Figure 2.5 displays E_z and E_r for both the driving laser pulse ($a_0 = 0.405$) and excited wake in the linear regime where $n_e = 1.5 \times 10^{18} \text{ cm}^{-3}$. Figure 2.6 displays E_z and E_r of a non-linear wake driven by an electron beam

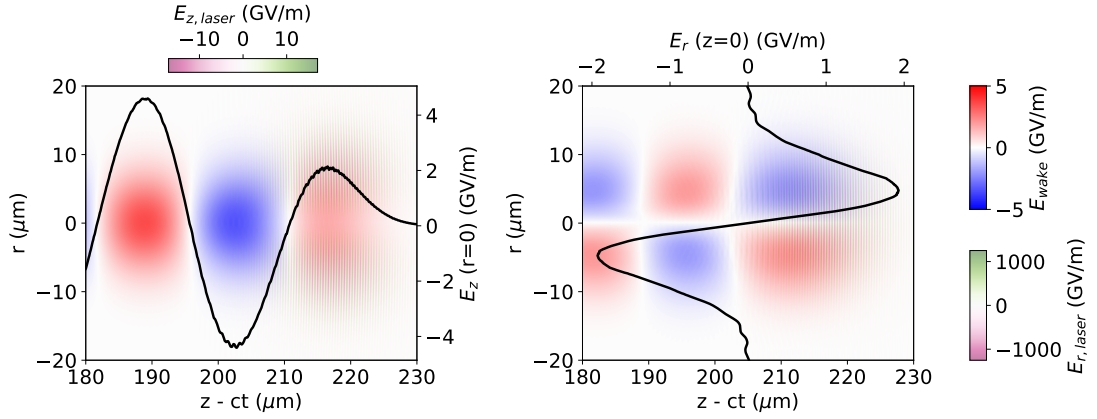


Figure 2.5: Linear regime of LWFA. The accelerating (left) and focusing (right) components of a wakefield generated in the linear regime ($a_0 = 0.405$) are illustrated. Electrons positioned within the accelerating phase of the nearly sinusoidal longitudinal electric field can gain energy from the wake. Just behind the driver, the transverse fields act to focus the electrons.

with $n_b = 6.6 \times 10^{18} \text{ cm}^{-3}$ where $n_e = 1 \times 10^{17} \text{ cm}^{-3}$. In the linear regime, accelerating and decelerating fields are symmetric in size and shape. Similarly, transverse focusing and defocusing fields are symmetric. As a result, there is approximately $\lambda_p/4$ phase region where electrons can be both accelerated and focused. In the non-linear regime, the drive beam almost completely evacuates the plasma electrons to create a region often referred to as the blowout or bubble. Notably, the accelerating and focusing regions for electrons are larger than those in the linear regime.

The maximum energy gain achievable over a fixed acceleration length is governed by the amplitude of the wake's accelerating field E_z . This amplitude is influenced by the density perturbations of plasma electrons within the density peaks $\delta n/n_0$. As the driver strength increases, the velocity of background electrons may surpass the phase velocity v_{ph} of the plasma wave, causing a disruption in the wakefield structure known as longitudinal wavebreaking. This phenomenon sets a maximum field limit E_{wb} beyond which the fluid plasma model is no longer valid. In the linear regime, E_{wb} can be estimated under the assumption that all

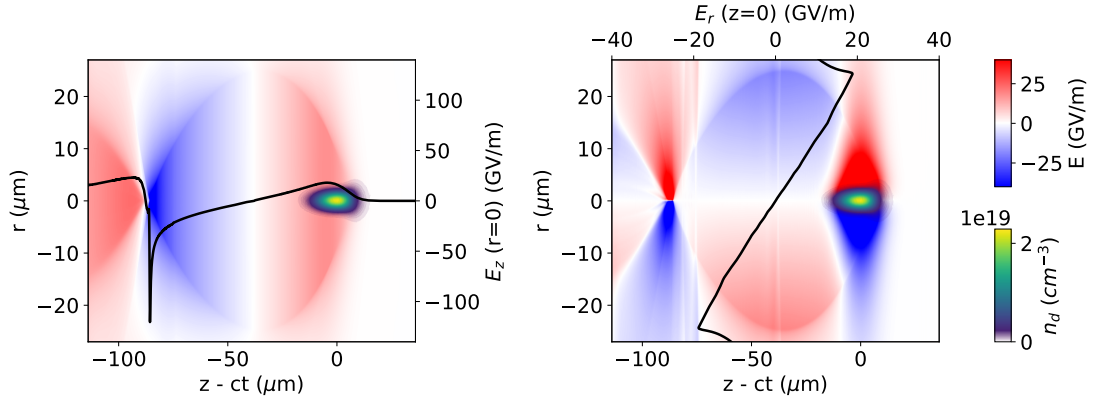


Figure 2.6: Non-linear regime of PWFA. The accelerating (left) and focusing (right) components of a wakefield generated in the non-linear regime ($n_b/n_e \simeq 66$) are illustrated. This regime can be identified by the distinctive sawtoothed shaped longitudinal fields and also focusing fields that remain linear across the entire blowout region, enabling the preservation of emittance in the accelerated electron bunches.

electrons of the wave undergo oscillations at ω_p . Solving Equation 2.1 under the cold, non-relativistic approximation yields the corresponding wavebreaking field

$$E_{wb} = \frac{m_e c \omega_p}{e} \approx 96 \sqrt{n_e \text{ (cm}^{-3})} \text{ Vm}^{-1} \quad (2.51)$$

In the non-linear regime, electron velocities approach relativistic magnitude, leading to appreciable Lorentz factor variations across different regions of the plasma wave. These differential phase velocities induce a phenomenon known as phase mixing, wherein adjacent electron fluid elements dephase and eventually intersect, resulting in wavebreaking. To incorporate these relativistic effects, the expression for the non-relativistic wavebreaking field is accordingly modified as follows

$$E_{wb,rel} = \sqrt{2(\gamma_p - 1)} E_{wb} \quad (2.52)$$

As in the non-relativistic case, background electrons can exceed the wave's phase velocity, leading to their injection into the accelerating phase of the wake.

Once initiated, this self-injection process typically persists, resulting in a broad energy spread. In three-dimensional geometries, the wave may also undergo transverse wavebreaking, which generally occurs at significantly lower driver amplitudes than predicted by the idealized one-dimensional model. To preserve high beam quality, it is often desirable to suppress self-trapping by operating below the wavebreaking threshold and employing controlled injection techniques instead (see Section 2.5.2).

2.5.2 Injection methods

Following the generation of a wakefield by a driver, the subsequent challenge lies in injecting a particle bunch into its accelerating phase. This process typically requires precise coupling of the beam into a narrow region of characteristic length k_p^{-1} , co-propagating with the wake at a phase velocity v_{ph} . The conventional trapping condition dictates that the injected particles must possess a velocity v_i satisfying $v_i \geq v_{ph}$ to remain within the accelerating phase of the wake. As outlined in Section 2.5.1, the wake phase velocity in PWFA approaches the speed of light ($v_{ph} \approx c$) thereby rendering the trapping condition considerably more strict than in LWFA, where v_{ph} is typically subluminal ($v_{ph} < c$). In general, injection techniques for introducing a witness beam into a plasma wake can be classified into two main categories: External injection and Internal injection.

External injection

External injection, particularly well-suited for PWFAs operating in low-density plasmas, involves the acceleration of pre-accelerated particles originating outside the plasma medium. In early PWFA experiments, electrons from the trailing edge of the drive bunch were accelerated [84]; however, optimal configurations employ a dedicated witness bunch, either sourced from an independent linac injector [85] or generated via controlled separation of a single bunch into driver and witness components [86]. A key advantage of external injection is that the injected

particles are typically already relativistic, inherently satisfying the trapping condition required for acceleration. The primary objective is thus to phase-match the witness bunch with the accelerating phase of the wakefield and sustain high-gradient acceleration while preserving critical beam parameters such as charge, energy spread, and emittance. Accelerator stages utilizing this injection scheme effectively function as energy transformers, converting energy from a high-charge driver into increased energy for a lower-charge witness beam.

Internal injection

Alternatively, witness bunches can be generated in situ within the plasma wake through internal injection mechanisms. In this scenario, injected particles typically possess subluminal initial velocities and must be rapidly accelerated to match the phase velocity of the wake in order to satisfy the trapping condition. This process is generally more accessible in LWFA due to the strong longitudinal electric fields and evolving wake dynamics, often to the extent that uncontrolled injection, commonly referred to as dark current, can occur [87]. In contrast, the more stringent trapping condition in PWFA makes internal injection more challenging, however this also enables the generation of witness beams with minimal or negligible dark current [88]. Various techniques have been proposed which exploit different mechanisms, however from an experimentalist's point of view, self-injection is seen as the most straightforward approach. It relies on the wave-breaking phenomena discussed in Section 2.5.1 and is often augmented by the expansion of the cavity size during for instance the self-focusing or self-compression of the drive laser in LWFA. However, witness beams created via this method tend to exhibit lower reproducibility, as they are inherently coupled to often random oscillations in the driving beam.

Another method for triggering internal injection involves the generation of electrons within the blowout region via field-induced ionization of background plasma ions or additional neutral dopant species. This ionization, typically oc-

curing through quantum tunneling (see Section 2.3.1), can be driven by the strong electric fields of the drive beam [89, 90], the wakefields themselves [91], or externally applied ionization lasers. The latter approach, comprising methods such as multi-pulse laser collisions [92] or the laser-triggered ionization of dopant gases - commonly referred to as the Trojan Horse or plasma photocathode technique [93] (see Chapter 4), offers the advantage of decoupling the injection trigger from the beam or wakefield dynamics, thereby allowing greater control and tunability. Electrons released near the axis in such ionization-based injection scenarios typically possess minimal initial transverse momentum, resulting in beams with ultra-low transverse emittance, an attribute often described as cold injection. Conversely, these electrons are born with negligible longitudinal momentum, which significantly increases the challenge of satisfying the trapping condition. As a result, these injection schemes are generally constrained to high-current driver beams (≥ 5 kA) operating in the non-linear, strongly blown-out regime, where the accelerating fields are sufficiently intense to capture particles born at rest.

A different yet increasingly common internal injection approach is to activate the self-injection method at a defined position along the propagation z by the controlled reduction of the wake's phase velocity v_{ph} in a negative plasma density gradient $\partial n_e / \partial z < 0$ [94–96]. In this so-called density downramp, the plasma wave elongates according to $\lambda_p \propto n_e^{-1/2}$ and the back of the plasma wave slows down momentarily, thus lowering the injection threshold and causing self-trapping. After the down-gradient, the drive beam propagates the wake in low density environment which accelerates the injected electrons, but avoids self-trapping. In a 1-D treatment, the wake phase can be approximated to $\chi = k_p(z)(z - ct)$ and thus the wake phase velocity in a plasma of density $n(z)$ is given by [97]

$$\frac{v_{ph}}{c} = -\frac{(\frac{\partial \chi}{\partial ct})}{(\frac{\partial \chi}{\partial z})} = \frac{1}{(1 + \frac{\xi}{k_p} \frac{dk_p}{dz})} \quad (2.53)$$

where $\frac{dk_p}{dz} = \frac{k_p}{2n} \frac{dn}{dz}$ and $\xi < 0$ behind the drive beam, highlighting the necessity of a negative density gradient. During the wake elongation, electrons originally located at a phase $\chi_i = k_{p,i}\xi$ will move and become located at $\chi_f = k_{p,f}\xi$, such that their wake phase advance is given by

$$\Delta\chi = \chi_f - \chi_i = \chi_i \left(1 - \sqrt{\frac{n_f}{n_i}}\right) \quad (2.54)$$

where n_i and n_f are the densities at the top and bottom of the downramp, whilst $k_{p,i}$ and $k_{p,f}$ are their corresponding plasma wavenumbers. The rephasing of plasma electrons into the accelerating region of the wakefield allows them to rapidly gain velocity, which when coupled with a reduction in the wake's phase velocity, can lead to electron trapping. Techniques for generating density downramps include the implementation of tailored gas jet profiles, the formation of hydrodynamic shock fronts [78, 80, 98], and the controlled optical ionization of higher-lying states in the background plasma or dopant species to create localized density spikes, commonly referred to as plasma torches [99]. While electrons are typically injected with residual momentum during such processes, it has been demonstrated that, under optimized conditions, density downramp injection can yield electron beams with phase space characteristics and overall beam quality comparable to those obtained via ionization injection [100].

Across the board, internal injection mechanisms in plasma accelerators hold the potential to generate particle beams with brightness surpassing that of conventional accelerator technologies. Moreover, if hybrid LPWFA plasma photocathode schemes, such as the one investigated in Chapter 4, produce witness bunches with brightness exceeding that of the drive bunch, they can effectively function as brightness transformers.

2.5.3 Wakefield acceleration

Once a particle bunch is successfully trapped within the accelerating phase of the wakefield, it can undergo rapid energy gain. The following section provides a brief overview of key metrics used to characterize acceleration quality, along with a discussion of relevant physical effects that may influence the acceleration process.

Transformer ratio

Where a plasma wakefield accelerator's primary aim is to function as an energy transformer, it is essential to maximize the efficiency of energy transfer from the drive bunch to the trailing bunch. To quantitatively assess this efficiency, the transformer ratio R [101] is introduced

$$R = \frac{E_+}{E_-} \quad (2.55)$$

where E_+ represents the peak accelerating field within the wake, preferentially coinciding with the field experienced by the witness bunch. In PWFA, E_- denotes the peak decelerating field acting on the drive bunch, where as in LWFA it refers to the maximum deceleration gradient experience by laser photons. E_- directly influences one of the key restrictions in PWFA, namely driver energy depletion, which can be defined as [102]

$$L_d = \frac{m_e c^2 \gamma_b}{e E_-} \quad (2.56)$$

where γ_b is the relativistic gamma factor of the drive beam. This effect arises when electrons in the drive beam are decelerated to sub-relativistic velocities, resulting in their entry into the accelerating phase of the wakefield. This backflow perturbs the wake structure and compromises the integrity of the acceleration process, thereby imposing an upper bound on the energy gain achievable by the trailing bunch. To mitigate this effect, it is desirable to maximize the peak ac-

celerating field E_+ while simultaneously delaying the onset of driver depletion by minimizing the peak decelerating field E_- . This can be accomplished by flattening the longitudinal decelerating field experienced by the drive bunch, ensuring that energy extraction occurs uniformly along its length [103]. Such field uniformity can be realised by modifying the current profile of the driver to conform to the spatial structure of the wakefield, thereby enabling ramped or triangular current profiles [104, 105] to optimize the value of R .

Depletion, Dephasing and Diffraction

Driver depletion is also an important limitation in LWFA, where it is often referred to as pump depletion. The characteristic length for pump depletion is defined by the distance in which the laser pulse loses half of its energy to the plasma wake, which in the 3D non-linear regime can be estimated as [106]

$$L_{pd} \simeq \frac{\omega_L^2}{\omega_p^2} c t_{FWHM} \quad (2.57)$$

Furthermore, as a consequence of the laser travelling at $v_d = c\sqrt{1 - \frac{\omega_p^2}{\omega_L^2}} < c$, electrons trapped in an LWFA wake can attain velocities closer to c , than the group velocity of the driving pulse. As a result, they gradually overtake the laser and undergo dephasing - that is, they transition from the accelerating phase into the decelerating phase of the plasma wakefield. This dephasing imposes another limit on the maximum energy gain achievable by the electrons, which is inherently linked to the relativistic plasma wavelength λ_p . The characteristic distance over which an electron remains in the accelerating phase i.e., before it encounters the zero-crossing of the longitudinal electric field ($E_z = 0$) in the laboratory frame, is referred to as the dephasing length, denoted L_{dp} , and is defined as [107]

$$L_{dp} = \frac{4}{3} \frac{\sqrt{a_0}}{k_p} \frac{\omega_L^2}{\omega_p^2} \quad (2.58)$$

In the linear regime, acceleration is generally limited by dephasing ($L_{dp} \ll$

L_{pd}), for example, when $a_0 = 0.5$, $\lambda_L = 0.8 \mu\text{m}$, and $n_e = 10^{18} \text{ cm}^{-3}$, $L_{dp} \sim 3 \text{ cm}$ and $L_{pd} \sim 45 \text{ cm}$. In the non-linear regime, L_{dp} increases while L_{pd} decreases, leading to a scenario where $L_{dp} \sim L_{pd}$. Both dephasing and depletion effects can be mitigated by reducing the plasma density n_e , which extends these characteristic lengths, however it is important to note that operating at lower plasma densities necessitates a correspondingly higher normalized vector potential a_0 and a longer laser pulse duration t_{FWHM} in order to sustain the equivalent level of wakefield excitation.

The third main limitation on LWFA is diffraction, a concept introduced earlier in this Chapter in Section 2.2, which reduces the laser intensity of any focused beam after a certain distance. However, either self-focusing, as discussed at length in Section 2.4.1, or a specifically engineered plasma channel with a parabolic density profile [108] can be invoked to counter-balance diffraction over several Rayleigh lengths.

Beam loading

Beam loading, defined as the perturbation of the accelerating field by the collective space-charge fields of the accelerated bunch, represents a fundamental effect in plasma wakefield accelerators, arising when the witness bunch induces its own wakefield which then interferes destructively with the wake driven by the driver beam [109]. This interaction becomes particularly significant for high-current trailing bunches, which can substantially deform the local longitudinal electric field, corresponding to an enhanced rate of energy transfer from the wake to the beam [97]. The concept of optimal beam loading is central to maximizing acceleration efficiency whilst preserving beam quality and can be brought about by tailoring the current profile of the witness bunch to match the spatial structure of the wake, typically approximated by a trapezoidal distribution [110], which in turn flattens the accelerating field along the bunch length. This field uniformity is critical for minimising correlated energy spread and preserving the longitu-

dinal phase space of the accelerated bunch, thereby enabling the generation of high-brightness beams.

2.6 Computational tools

While the theoretical framework presented in this chapter to date is essential for understanding the underlying physics of plasma-based accelerators, precise replication of many critical phenomena, such as the highly non-linear dynamics of the blowout regime, requires the use of numerical simulation techniques. A widely adopted numerical technique for simulating plasma accelerators is the Particle-In-Cell (PIC) method [111]. This approach leverages the collective behaviour of plasmas by discretizing both space and time: rather than solving the governing equations over a continuous spatial domain, the simulation domain is divided into a finite number of spatial cells, and the system is evolved in discrete time steps Δt

To reduce computational complexity, the PIC method does not track individual electrons and ions. Instead, it represents large ensembles of physical particles as macroparticles. Each macroparticle carries the statistical average of the position and momentum of many real particles, while maintaining the correct charge-to-mass ratio (q/m) through assigned particle weights. Preserving this ratio ensures that macroparticles respond to electromagnetic fields according to 2.14 in the same manner as individual particles, allowing the dynamics to remain physically accurate. This provides a computationally efficient yet self-consistent procedure for simulating the complex plasma phenomena encountered in plasma-based acceleration.

Figure 2.7 illustrates the general PIC algorithm cycle, highlighting the fundamental steps that are common across the various PIC codes [83, 112–114] used in plasma accelerator modelling. The iterative cycle is as follows:

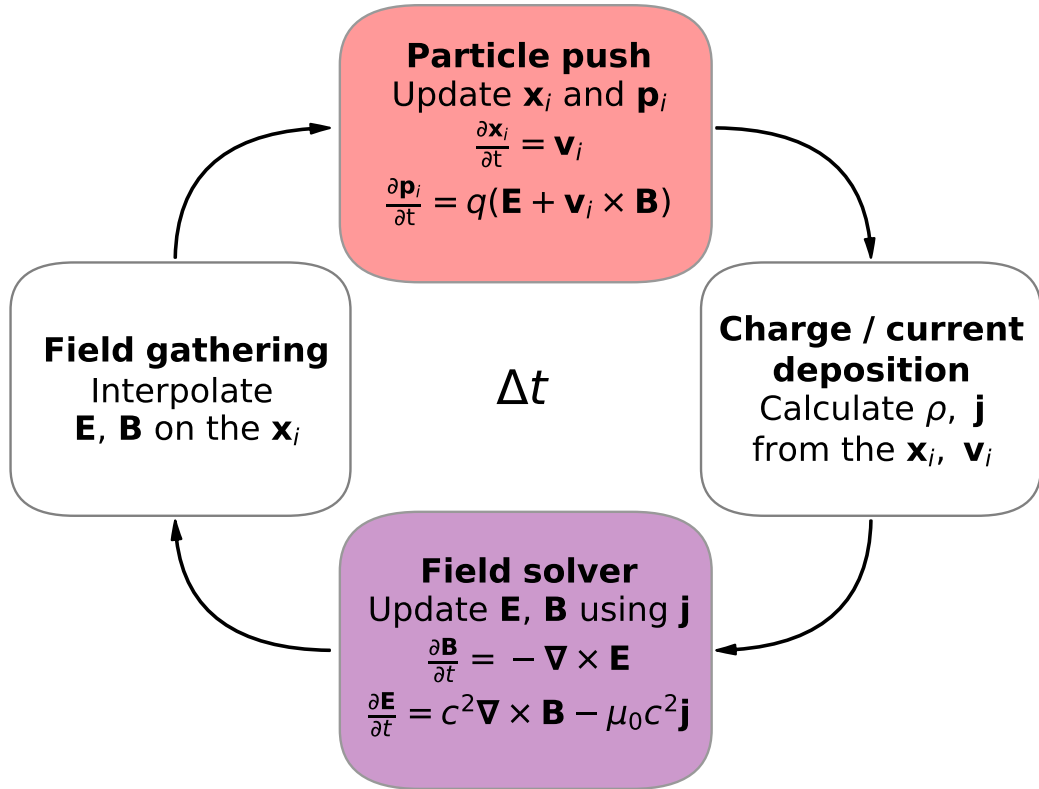


Figure 2.7: Generalised PIC cycle. At each time step, the simulation performs a full cycle of operations. The positions $\mathbf{x}$ and velocities $\mathbf{v}$ of the macro-particles are first used to compute the charge density ρ and current density $\mathbf{j}$. These quantities are then used, via Maxwell's equations, to calculate the updated electric $\mathbf{E}$ and magnetic $\mathbf{B}$ fields using the leapfrog integration method. Finally, the particle pusher updates the motion of the macro-particles by applying the Lorentz force, and the cycle proceeds to the next time step.

1. **Particle push:** The macroparticles are advanced in time using discretized forms of the Newton–Lorentz equations of motion, based on the electromagnetic fields computed at each time step.
2. **Charge/current deposition:** Based on the particles’ positions and momenta, the corresponding charge and current densities are computed and deposited onto the simulation grid.
3. **Field solver:** Maxwell’s equations are solved to obtain the electromagnetic fields generated by the computed charge and current distributions.
4. **Field gathering:** Electromagnetic fields are interpolated to the macroparticle positions, typically using a linear interpolation scheme such as the nearest-neighbour method.

Following step 2, once ρ and $\mathbf{j}$ have been deposited onto the grid, additional physics modules such as field-induced ionization models [31, 115] (see Section 2.3.1) are typically evaluated at this stage. These enable the self-consistent generation of new particle species in response to the local electromagnetic fields. While beyond the scope of this work, other physical processes, such as collisional effects, can also be incorporated into the PIC cycle through appropriate extensions. Consequently, the careful selection of a PIC code that supports the required physical models and aligns with the objectives of the simulation study is essential. As with any numerical simulation, resolution plays a critical role in the accuracy of PIC modelling. Both the spatial cell size and the temporal step size must be sufficiently small to resolve the relevant physical phenomena. A standard approach to determining adequate resolution involves progressively reducing the cell size and/or time step until key output parameters exhibit negligible variation. When further refinement yields no significant change in results, the simulation is considered to have achieved convergence.

The simulation work presented in this thesis utilised two distinct PIC codes. During the initial exploratory phase of the study detailed in Chapter 4, simulations were conducted using VSIM [112], a general-purpose, fully 3D Cartesian PIC code. VSIM has a well-established track record in PWFA research [47, 116–118], with one of its primary strengths being the modularity of its input deck structure. This modularity, combined with its flexibility in accommodating arbitrary particle and field geometries, makes it particularly well-suited for the wide range of configurations typically encountered in PWFA studies. However, like all Maxwell solvers based on the Finite-Difference Time-Domain (FDTD) method and the Yee grid [119], VSIM is susceptible to Numerical Cherenkov (NC) instability. This numerical artifact arises from the coupling of spurious high-frequency electromagnetic modes, typically at frequencies ω_{NC} , to particles propagating at subluminal velocities. Although this effect can be partially mitigated by selecting a time step that satisfies the Courant–Friedrichs–Lewy (CFL) stability condition [120]

$$\frac{c\Delta t}{\sqrt{(\delta x)^2 + (\delta y)^2 + (\delta z)^2}} < 1 \quad (2.59)$$

its complete suppression often necessitates significantly finer spatial resolution. This, in turn, results in substantially increased computational cost, rendering such approaches impractical for many large-scale or long-duration simulations.

For the production runs presented in Chapter 4, as well as all simulations discussed in Chapters 3 and 5, FBPIC [83] code was employed. Designed specifically for plasma accelerator applications, FBPIC operates in a quasi-3D cylindrical geometry, enabling significantly higher resolution at a reduced computational cost compared to fully 3D Cartesian codes. In contrast to VSIM, FBPIC is a relatively recent but increasingly established simulation framework for modelling both laser and beam-driven plasma acceleration [72, 77, 80, 121, 122]. It utilizes a spectral Maxwell solver, which is inherently immune to NC instability, eliminating the

need for overly restrictive time steps or excessive spatial resolution. Additionally, its GPU-accelerated architecture facilitates orders-of-magnitude faster computation than traditional CPU-based 3D PIC codes. While particle trajectories are tracked in full three-dimensional space, the electromagnetic fields are reconstructed using azimuthal mode decomposition, making FBPIC particularly well suited to quasi-axisymmetric problems, such as conventional LWFA or plasma filamentation.

Chapter 3

Kilohertz-rate laser-plasma acceleration

As discussed in Chapter 2, Laser-plasma acceleration (LPA) is an innovative technique for accelerating charged particles to extremely high energies over a very short distance [58]. It has demonstrated the ability to generate accelerated electron bunches with both femtosecond pulse duration and high peak current, making them attractive for a number of scientific and industrial applications [123–125].

At present, 100 TW or even PW-scale laser systems are typically employed to operate such accelerators. Utilising lasers with such high input power allow laser-plasma accelerators to pursue electrons with energies in the GeV range, capable of driving advanced light sources such as free electron lasers or examining the fundamental forces of matter in linear colliders in high-energy physics [126]. However, these laser systems are up to date only available in dedicated research facilities and their operation is very demanding, requiring large, highly specialized groups with strong expertise. They are also restricted to run at a few Hertz repetition rate or less which limits their practical use [127].

Therefore, increasing LPA drive laser repetition rates to the kHz range is

beneficial not just to expand the technique’s usability, but also for a host of other reasons. Firstly, from an experimental perspective, it will improve the data acquisition process by enabling rapid averaging over many shots, thus increasing a measurement’s signal to noise ratio. Secondly, from a physics perspective, it allows for the laser to reach a thermal steady state leading to increased laser and subsequent electron beam stability [128]. Thirdly, it opens the door for active feedback control of the experimental setup to optimize the output electron beam’s parameters in real time [129]. Finally, the electron current/flux can be magnified by orders of magnitude, further broadening the applicability of LPA’s.

Currently available kHz systems are smaller both in terms of size and power, typically high GW-scale, than traditional LPA lasers. However, despite providing less energy, they have already proven their potential as few MeV laser–plasma accelerators [130] and for delivering e-beams with long-term stability [131]. Electrons with energies in the 1-10 MeV range are in fact ideally suited for several applications such as pulsed radiolysis [132] or ultrafast electron diffraction (UED) [133].

This chapter lays out advancements made towards laser-plasma acceleration at high repetition rates on two separate fronts. Firstly, as part of a collaboration with academic partners and the company SourceLAB, the author contributed towards the design, testing and verification of e-KAIO (electron Kit-All-in-One), a laser–plasma electron accelerator module designed to be driven specifically by high repetition-rate lasers for industrial applications [16]. This marks a first step towards making LPA technology more accessible to a broader community in scientific research and industry. Secondly, the author conducted a simulations based study to demonstrate the ability of contemporary kHz lasers to create electron flux with attributes akin to those found in space radiation environments, specifically the outer Van Allen belt, which can subsequently be used to satisfy the growing demand for suitable beam sources to perform radiation hardness

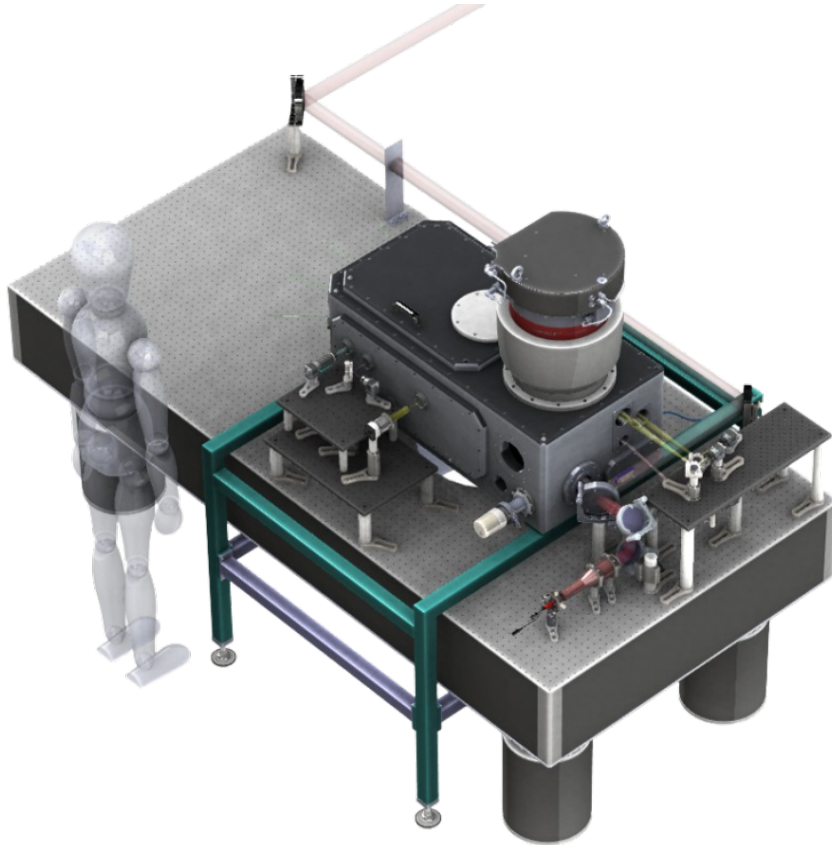


Figure 3.1: 3D e-KAIO visualisation. Overview of the e-KAIO acceleration module on an optical table. The compact design allows to employ the chamber for various applications.

assurance testing of electrical components such as satellite electronics [134].

3.1 e-KAIO

3.1.1 Description and design considerations

Despite a growing desire throughout the field to move towards higher repetition rates, implementing kHz LPA sources remains challenging as they require high-performance vacuum handling to manage the large target gas loads, and adequate management of large data sets generated during measurement scans.

The novel e-KAIO is designed to be a user friendly solution to these challenges allowing for fine adjustment of the parameters essential for high-quality beams such as control of the laser beam alignment, control of focal spot quality, control of gas jet position, density and profile, and the electron spectrometer dispersion. It consists of a single vacuum chamber containing all the necessary components for producing, optimizing, and monitoring electron beams generated via laser plasma acceleration in a gas jet, coupled to an automatic high repetition rate data acquisition program. Furthermore, e-KAIO is highly versatile and can be used in combination with different driving lasers for various applications.

The vacuum chamber is shown in Fig 3.1 and can be operated even in small experimental areas. The vacuum chamber itself is carried by an outer support structure, and the inner breadboard is mechanically isolated from the chamber walls by flexible bellows to reduce vibrations produced by the pumping system or the surroundings. The dimensions of the vacuum chamber itself are 500 mm $\times$ 900 mm $\times$ 450 mm, and the optical entrance window supports beam diameters up to 50 mm.

The e-KAIO chamber provides multiple diagnostics for monitoring the laser beam or studying the electron beam. The optical diagnostics are installed outside of the chamber including six CCD cameras that are connected to the KAIO control computer. A detailed overview is shown in Fig 3.2. The input beam is recorded by two reference CCDs which can be used for alignment references, thus enabling a quick realignment procedure from day-to-day operation. In addition, the references can be used for active (automatized) beam stabilization as they are coupled to piezo-motors controlling the position and pointing of the input beam. In this mode, the drift or misalignment of the laser beam will be corrected automatically by moving the pico-motors of the two last mirrors before the focusing off-axis parabola. This option is useful during the pumping of the chamber under conditions of strong vibrations and/or mechanical stress, as well

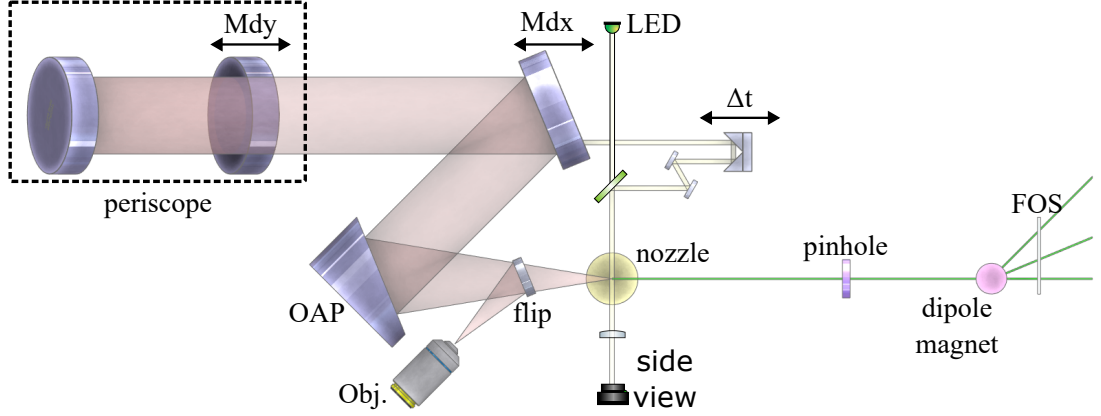


Figure 3.2: Schematic of the e-KAIO acceleration module. After characterization, a suitable laser beam is expanded by a mirror telescope and coupled into the e-KAIO chamber. The beam is reflected to the upper level by a 45° periscope of mirror Mdy and the upper turning mirror. The mirror Mdx reflects the beam to the off-axis parabola OAP, which focuses the beam into the gas jet above the nozzle. The gas jet is imaged from the top (not displayed) and the side to the respective cameras CCD TopView and CCD SideView (see the GUI in Figure 3) outside the chamber. Additional laser-plasma diagnostics could be added in either arm (outside the vacuum chamber) to investigate the light emitted from the plasma. The flip mirror in between the OAP and the jet can be used to deflect the attenuated beam to microscope objective Obj., which in turn images the focus onto a camera outside the chamber. The interaction volume can be illuminated from the side either by a continuous white light LED or a small sample beam, extracted from the main beam by a small hole on the side of the mirror Mdx. The quality of the focal spot can be optimized by iteratively adjusting the tip and tilt angles of Mdx and the OAP and the beam pointing is corrected by performing a translational shift before focusing by translating the mirrors Mdy (up/down) and Mdx (left/right). The translational shift leads to a rotation of the laser axis around the focal point and thus can be used to align the laser along the acceleration axis (green). Once aligned, the components of the electron spectrometer (pinhole, dipole magnet, and FOS) can be positioned by moving their respective translation stage.

as during long-term system operation. The alignment references are shown in the upper left of Fig 3.3, and the integrated signal of one of the cameras can be used to track the average power stability of the laser. In addition, the deviation from the set reference can be recorded to investigate the beam drift and pointing under different conditions. For pre-alignment purposes, the e-KAIO chamber is equipped with a red alignment laser diode, which has the same path and close central wavelength as the actual kHz laser. This allows one to align equipment without having to use a high-intensity laser. The mirrors and off-axis parabola used for focusing the beam are motorized and can be positioned via the control unit and the dedicated computer. The beam is focused by a f/3 or f/6 off-axis-parabola and the attenuated focal spot can be imaged by an objective when the flip mirror in front of the gas jet is up.

The gas jet is installed on a 3-directional translation stage and can be moved with micrometre precision. During the interaction of the high-intensity laser beam and the gas jet, electrons are accelerated which follow the axis of the main laser beam direction. These electrons are detected via a scintillator (Hamamatsu FOS J6677) screen, and the converted visible light is further recorded by a CCD camera after passing an interference filter.

Additionally, a magnetic spectrometer can be implemented in the electron beam axis as needed. The spectrometer consists of a pinhole and a pair of permanent magnets positioned on individual translation stages. By adjusting the distance of the magnets, the magnetic field strength can be adjusted to vary the electron dispersion. The field strength is pre-calibrated, and the dispersion axis is shown in the user interface (upper right in Fig 3.3). At the back of the interaction chamber are multiple vacuum flanges as a feed-through for additional diagnostics, to connect a beamline for further manipulation of the electron beam or to measure the source output (such as beam charge).

The e-KAIO also has a Graphic User Interface (GUI) that allows online control

of all parameters and data acquisition. In Fig 3.3, an overview of the e-KAIO chamber’s GUI is shown. The interface supports the multiple CCD cameras, the several motorized translation stages, the piezo-motors for driving the mirrors, and also the data acquisition and storing. The interface not only allows for monitoring and adjusting the parameters, but it also contains quick analysis of the focal spot and electron density for instance. Moreover, the e-KAIO software (version 1.0) has an implemented GS algorithm to measure the spectrally averaged wavefront of the laser beam. The recorded images and values are stored in HDF5 files giving the possibility to store large amounts of data. To reproduce and keep track of the experimental conditions used, a tagging function is included in the HDF5 file-saving option. Finally, a tailored e-KAIO HDF5 viewer was constructed to provide a more accessible user interface to the recorded data. This code was written predominantly using the Python package, PyQt5, and it permits the user the ability to instantaneously visualise any of the captured parameters by simply interacting with the self-explained pop-up window.

3.1.2 Experimental testing and verification

The first experimental efforts to generate electrons with a kHz-laser system employing the e-KAIO chamber made successful progress in bench-marking the system. Various diagnostics and subsystems were tested and proven to be operational.

The basic e-KAIO vacuum pumping system is a combination of one primary rough pump (Edwards iXL120, West Sussex, UK) and two turbo pumps (Pfeiffer HighPace[®] 2300, Asslar, Germany) with a suction of 2000 L s^{-1} each. In addition, a differential pumping configuration can be implemented to reduce the load on the turbo pumps. However, for these first proof of principle experiments, the e-KAIO was equipped with a pumping system consisting of two roughing pumps with a suction power of $120 \text{ m}^3 \text{ h}^{-1}$, one for the differential pumping close to the nozzle

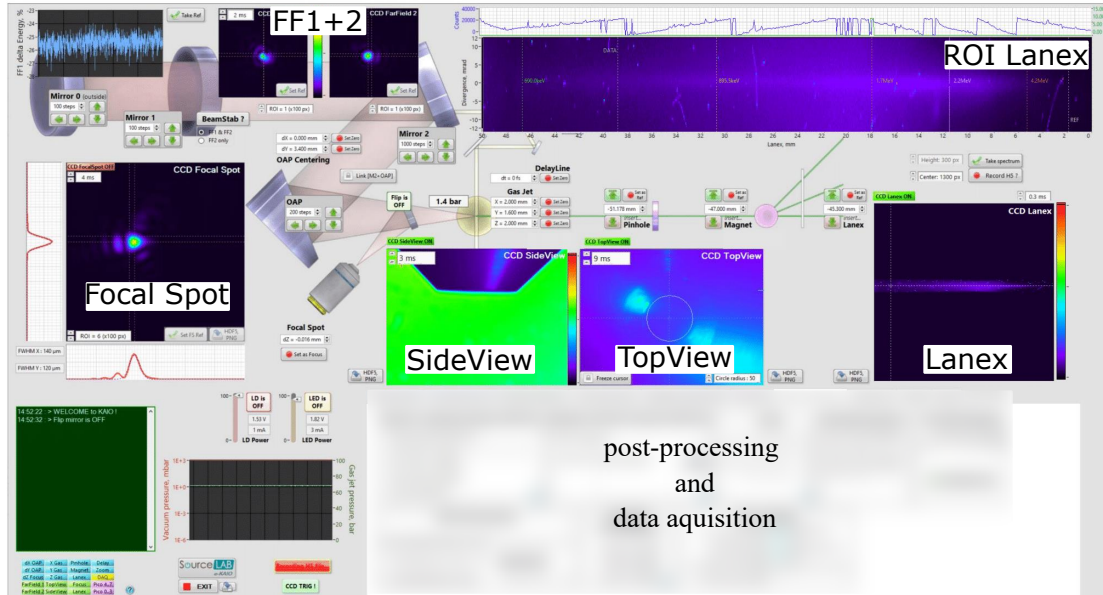


Figure 3.3: e-KAIO graphical user interface (GUI). An overview of all CCD cameras and motorized stages is given. In the upper left, the laser input power is monitored, and the two far-field references of FF1 and FF2 are displayed. Below, the magnified far-field diagnostic CCD FocalSpot is shown when the flip mirror is activated. In the centre of the GUI, the SideView (spatial resolution $2 \mu\text{m}$) and TopView images of the gas jet are shown. The raw image of the CCD Lanex is displayed on the far right, and the postprocessed image ROI Lanex is on top in the upper right. The calibrated energy axis for the given magnet distance is set below the graph, and a line-out is shown above the graph. The range of interest (ROI) can be chosen in advance. The bottom left shows a status overview of the different components. Next to it, the alignment laser and nozzle illumination can be activated. A more detailed overview can be found in the SourceLAB user manual.

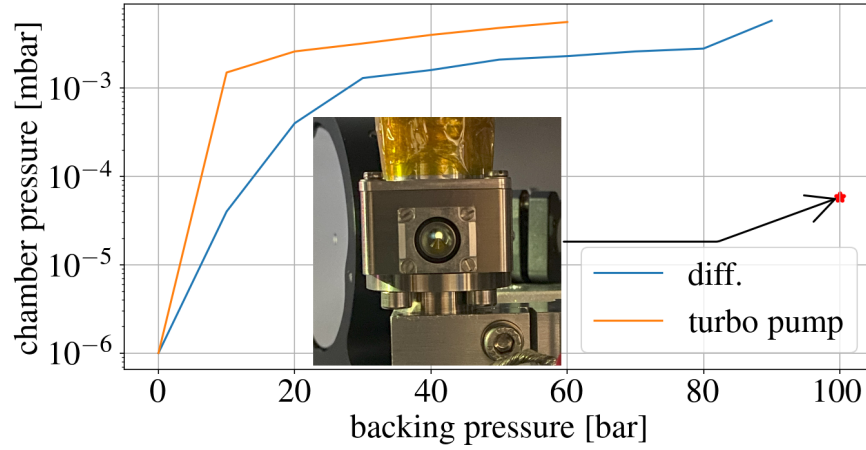


Figure 3.4: Pumping system gas load measurements. Pressure measurements for different pumping configurations. The chamber pressure is measured for different backing pressures of the nozzle for three different cases. The orange curve shows the pressure for the basic pumping (only turbo pump), and the blue curve presents the pressure for the rough differential pumping. A more promising differential pumping system was realized with a small cube attached to the valve. The pressure was measured only for a single backing pressure, which is indicated with the red star. The insert shows the differential pumping cube attached to the valve. The cube is connected to the rough pump by a flexible vacuum tube, and the side view diagnostic is maintained by two thin fused silica windows.

and the second one as a backing pump for the upper turbo pump. The connection of the optional bottom turbopump was covered by a blind flange. A pressure gauge directly connected to the e-KAIO control unit enables monitoring and saving of the values continuously via the KAIO GUI. Without gas load, chamber pressures of low 10^{-6} mbar to a few times 10^{-7} mbar could be reached. Fig 3.4 shows the measurements for the different pumping systems. In the differential pumping case (blue line), a tube connected to a roughing pump was placed a few centimetres on top of the gas nozzle to reduce the load on the turbopump. This simple differential pumping system already improved the chamber pressure enough to be able to operate the nozzle at a backing pressure of 90 bar. Without a differential pumping system (orange line), a chamber pressure of 5.8×10^{-3} mbar is reached at a backing pressure of 60 bar due to the high load on the turbo pump. To further optimize the gas load handling, a more advanced differential pumping scheme was implemented, where the interaction region is enclosed by a differential pumping cube. Thus, the gas leak into the vacuum chamber is limited to the openings for the laser beam. The probing of the interaction region is maintained by thin windows on either side of the cube. The valve of the gas nozzle is driven via PulseLAB, a digital pulse driver manufactured by SourceLAB. The device is triggered externally and can for instance be synchronised with the laser pulse. PulseLAB allows different modes to operate the nozzle: single-shot, continuous, or burst mode. The delay and opening time of the valve relative to the external trigger can be set in the control panel or in the PulseLAB software.

Two different gas nozzles were used to test the diagnostics and gather the first experimental results. A metallic symmetrically shocked nozzle [135] provided by SourceLAB and a supersonic glass nozzle produced by LightFab were tested using argon as a target gas. Due to manufacturing reasons, a glass nozzle could provide a smoother density profile. A drawback of the glass nozzles is their lower robustness to high backing pressure. The gas density was determined by a SID-4HR

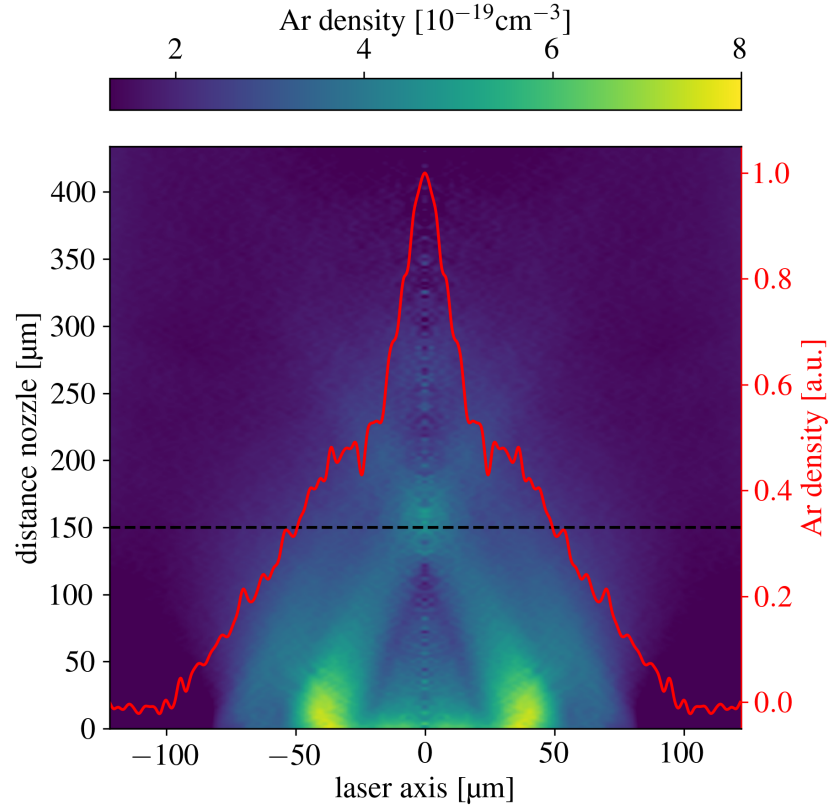


Figure 3.5: PHASICS gas profile lineout. Two-dimensional density distribution of a symmetrically shocked gas jet. The density profile (red) is taken at a height of 150 μm above the gas jet, where the shock fronts converge. The FWHM is around 50 μm and the spatial resolution is 2 μm .

PHASICS camera [136, 137] for interferometric measurements placed in the side view arm of the e-KAIO chamber (see Figure 3.2). Using the PHASICS camera software (version 6.3.1) add-on density module (not shown in the GUI) and the white backlight of the KAIO chamber, the target (atomic and plasma) density profile can be determined. The module performs an inverse Abel transformation of the phase shift projection to extract the atomic gas density by assuming a rotational symmetric gas jet. This is well justified for round nozzles. The resulting atomic density was found to be $3 \times 10^{19} \text{ cm}^{-3}$ for a backing pressure of 100 bar. For a fully ionized argon plasma ($Z = 8$), we assume a plasma density of $n_e = 2.4 \times 10^{20} \text{ cm}^{-3}$. The measured atomic density is presented in Figure 3.5. The FWHM of the gas jet profile is about $50 \mu\text{m}$. For imaging the gas jet onto the camera, a lens with a focal length of 10 cm was placed on the axis of the side-view diagnostics. The spatial resolution of $2 \mu\text{m}$ is limited by the imaging setup (magnification factor 15) and pixel size of the PHASICS CCD.

In the first phase of the experiments, a Coherent Astrella Ti:Sa laser system was used to drive the electron source. The Astrella laser system delivers up to 9 mJ/pulse and a pulse range of 35–100 fs at a repetition rate of 1 kHz. The nominal beam diameter is around $w \approx 12 \text{ mm}$ at a divergence angle of $\theta \approx 1\text{--}2 \text{ mrad}$. To achieve a sufficiently small focal spot size, the Astrella beam is expanded by a factor 4 mirror telescope to a diameter of around $2w \approx 50 \text{ mm}$ (maximum free aperture for the current design) before the focusing optic. The focal spot of the beam focused by an $f/3$ parabolic off-axis mirror is shown in Figure 3.6. The graphs show a line-out in the vertical (red) and horizontal (green) directions. The focal spot size (FWHM) is found to be approximately $4.00 \mu\text{m}$ in the horizontal direction and $4.70 \mu\text{m}$ in the vertical direction. The focus has a trefoil aberration, which is due to the manufacturing limits of the metallic parabolic mirror. This aberration leads to an energy loss of the laser focus in the wings of the focus (Strehl ratio $S \approx 0.8$). Considering the additional losses by the optics between

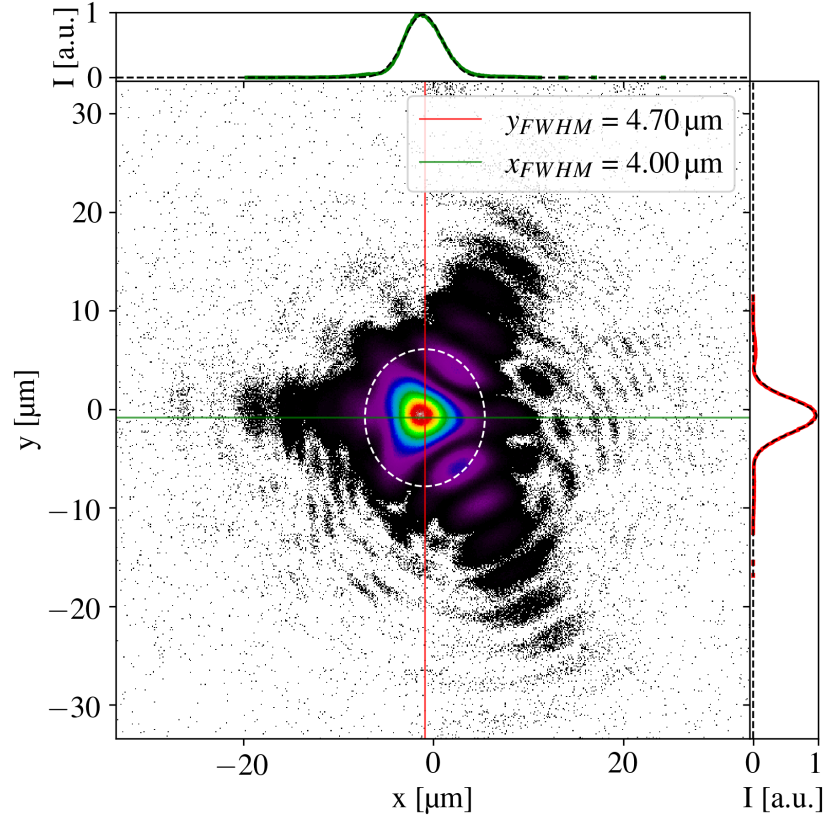


Figure 3.6: Coherent Astrella far-field beam profile. Focal spot intensity distribution recorded by the KAIO software. The centroid is determined, and the white ellipse marks the $D4\sigma$ region. The horizontal and vertical line-outs are plotted on the top and side, and the FWHM size is determined by fitting a Gaussian function. The trefoil aberrations are induced by the OAP. The spatial resolution of the imaging system is $0.1 \mu\text{m}$ (magnification factor 31).

the laser exit and the focus and the spread in time due to dispersion effects, the main pulse has an estimated peak intensity of $I_{peak} \approx 4 \times 10^{17} \text{ W/cm}^2$. This calculation is based on the pulse energy of 5.5 mJ measured on target and a pulse duration of 44.5 fs as found by the TIPTOE measurement. The high loss in pulse energy is due to numerous reflections from broadband silver mirrors and a slight clipping after beam expansion. The laser-plasma interaction was visualized via a shadowgraphy imaging line. The side-view imaging line is equipped with a delay stage to capture different moments of the laser-plasma interaction. The exact position of the mirror to adjust the delay Δt can be controlled within the GUI. Figure 3.7 shows the interaction at two different delays. The left-hand-side image serves as a reference since in this case no laser pulse was present. In the right-hand side image, the propagation of the laser pulse through the gas target can be observed. The delay between the laser pulse and the probe pulse on the right side of Figure 3.7 is 17 ps. The delay is marked in the top-right corner of each image. The Airy pattern in the background is due to the probe beam passing through the holed mirror when extracted from the main beam. An interference filter selecting narrow-band wavelengths around 800 nm is placed in front of the CCD camera to block the plasma light and other residual ambient photons.

3.1.3 Summary

To conclude, these first proof-of-principles experiments demonstrated the capability of the e-KAIO chamber to accommodate an industrialized laser system, and show that alignment and data acquisition using the e-KAIO chamber is straightforward and can be easily carried out in different laboratories or working environments. The first line of Table 3.1 summarizes the experimental results of this paper. Unfortunately, we were unable to accelerate electrons to significant energies, mainly because the laser’s focal intensity (or normalized vector potential a_0) was limited by the OAP’s aberrations and the pulse’s third-order

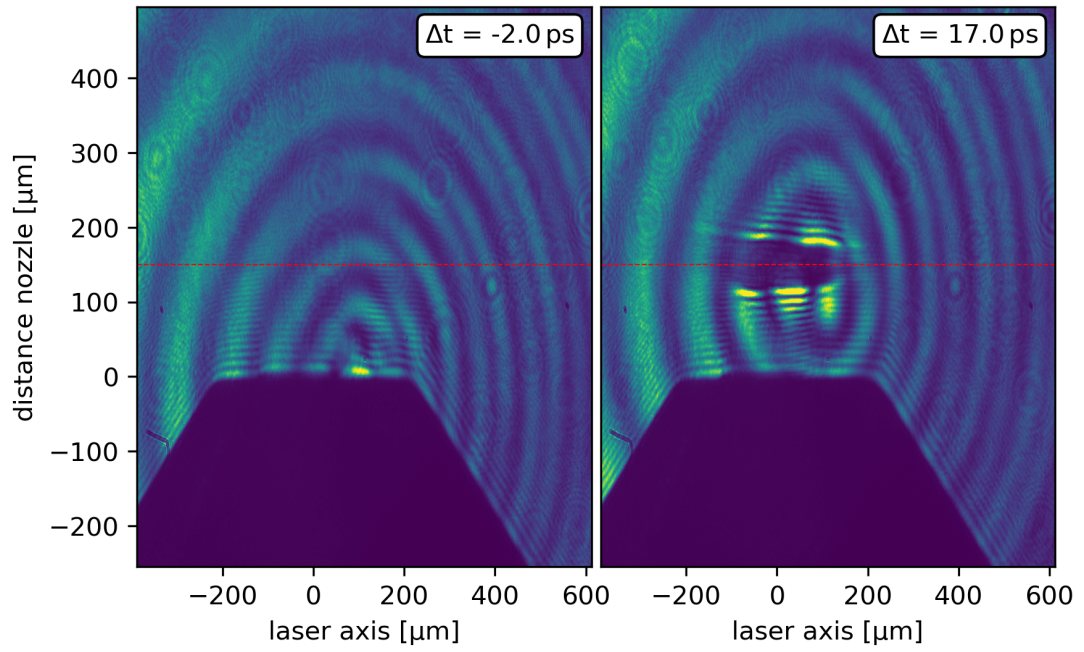


Figure 3.7: Side view shadowgraphy comparison. Shadowgraphy images of the interaction of laser pulse and target gas. The (left image) shows the reference (2 ps before the laser pulse arrives) and the (right one) the interaction at a delay of 17 ps. The red line marks a height of 150 μm , where the shock fronts of the nozzle converge.

dispersion. This then restricts self-focusing, as we are unable to reach the relativistic self-focusing threshold in a plasma ($P/P_c \approx 1$) at the measured target pressures, which is essential for electron injection in this case.

E_L (mJ)	τ (fs)	a_0	n_e (n_c)	P (P_c)	w_0 (μm)	z_R (μm)	ΔE (MeV)	acc. Regime
5.5	44	0.5	0.14	1	4.4	54	-	-
5	7	1.2	0.14	5.8	4	50	3	bubble
5.5	30	0.6	0.7	7.4	4	50	0.1	self-inj.

Table 3.1: Experimental parameters and their associated acceleration regimes. Comparison of different experimental parameters and their potential electron acceleration regimes. The first line highlights the results presented in this paper, the second line represents potential experiments in the bubble regime when the pulses are postcompressed to 7 fs, and the third line considers a potential experiment using a high-pressure gas jet (five-fold increase in atomic density). The two last lines include the assumption of aberration-free focusing optics and an optimized laser pulse dispersion.

A first step toward low-energy electron acceleration (keV level) working in the self injection regime would be the optimization of the focal spot and the pulse duration down to 30 fs, while increasing the target pressure to trigger self-focusing at these higher gas densities. A further increase in the plasma density could be reached by using molecular nitrogen as the target gas (10 electrons per molecule). Potential experimental parameters are given in the third line of Table 3.1. To reach electron energies in the MeV range with this experimental setup, a post-compression scheme needs to be implemented to reach the bubble or resonant regime [23] (second line in Table 3.1). A multi-pass-cell post-compression scheme with a total transmission of around 80% is planned for future experiments.

The e-KAIO interaction chamber is a first step in reaching the goal of a table-top laser-plasma electron accelerator on a university scale accessible for different

potential applications. It can provide a structural change in the availability of laser-plasma accelerators that will advance the research and progress of developing a system suitable for industrial applications. This will give opportunities to drive the research forward and make electron accelerators available on a smaller scale. Therefore, research and applications will advance further.

3.2 Space radiation reproduction

Over the next century, the frequency of space expeditions is set to increase globally with planning already underway for several missions beyond Low Earth Orbit (LEO) [138]. In 2017, NASA’s Artemis program was launched with the initial objective of installing a viable habitation on the Moon, and long-term aspirations of dispatching humans to Mars [139]. The radiation environment that will be encountered on such missions is highly complex and will likely undergo substantial changes during the course of the mission.

A particular area of interest is the Van Allen Belts [140], which are comprised of large fluxes of ionizing particles trapped and retained around the Earth by its magnetic field. Within the Outer belt, the energy of certain electrons can rapidly surge from the keV range to the MeV range [141, 142]. These relativistic electrons have been aptly labelled ‘killer electrons’, after research confirmed their ability to induce harmful radiation damage to both astronauts and spacecraft electronics [143]. For instance, there are several high-profile examples, including Hipparcos and ETS-6 satellites, of space probes that have developed faults as a result of interacting with these high energy particles [144].

Electronic components and circuits are essential to the success of space missions and satellites, with ever increasing demands being placed on their performance. Consequently, there is a growing need for beam time and more suitable radiation sources to test the effect these ‘killer electrons’, and other abstract types

of radiation, may have on proposed electronics. However, as previously discussed in Chapter 1, RF accelerators are limited in number due to the large costs associated with building and maintaining them. There is also a further issue in that conventional accelerators such as linacs and cyclotrons are designed to produce e-beams with nearly monoenergetic distributions. On the other hand, space radiation in the form of ‘killer electrons’ is very broadband and can be represented by an exponential or power-law energy spectra [145].

As such, Laser-plasma accelerators have been proposed as a more appropriate alternative for space radiation testing [146, 147] given that broadband radiation is the inherent regime of LPAs, with monoenergetic beams being more difficult to generate. They are also far more widely available, compact and cost effective than other traditional radiation sources. Recent proof-of-concept experiments have demonstrated the ability of LPA’s to reproduce the spectral characteristics of space radiation and shape the particle flux realised by tweaking the LPA interaction parameters [148].

To date, LPAs utilised for space radiation reproduction have relied upon focusing a TW-scale laser pulse onto a thin metal foil target. This laser-overdense plasma interaction is an incredibly effective and reliable method for converting laser energy into broadband electron flux, and also into protons via the Target Normal Sheath Acceleration (TNSA) mechanism [149, 150].

However, these accelerators have historically been limited to repetition rates of below 1 Hz because of the need to replenish the target foil, which is destroyed during the acceleration process, resulting in a low average flux. Despite significant progress in the development of tape-drive targets that now allow for operation up to the multi-Hz level [151], kHz scale repetition rates are required to achieve average flux levels that will allow for ‘accelerated lifetime testing’, such that the dose received by a satellite over the course of say 20 years can be delivered to the sample in a manageable time frame.

Advantageously, the combination of a gaseous jet target, that can achieve a near continuous gas flow, and a high rep rate laser used to drive a laser-underdense plasma interaction, may create a beam source capable of operating at repetition rates in the order of kHz or higher, increasing the output particle fluence/flux and satisfying the conditions required to mimic the total radiation dose accumulated by a spacecraft’s electrical component over its operational lifetime. This new setup could also facilitate higher peak fluxes (close to target) if required, which would be useful for studying collective or non-linear radiation effects.

Additionally, higher electron energy ranges (10-100 MeV) are more accessible via underdense, gaseous targets when compared with overdense, metallic targets. This is of particular relevance for modelling the radiation belts of other magnetised planets such as Jupiter, that are populated with energetic electrons of up to 150 MeV [152].

In this subchapter, a simulations-based study was performed using the PIC code FBPIC to model the aforementioned laser-underdense plasma interaction and identify a regime capable of producing the relativistic, broadband electron beams with energies up to a few MeV that are located within the Earth’s Outer Van Allen belt. In order to achieve this, the challenges and features of operating plasma accelerators using current kilohertz laser systems such as self-focusing, self-compression and high-density plasma sources were explored so they could be better understood and exploited to achieve the desired output electron spectra. Finally, once this acceleration regime was identified, fluctuations to key experimental parameters (laser energy, pulse duration, gas density etc.) were induced to allow for the overall stability of the accelerator to be diagnosed. This was followed by a discussion of potential methods to improve the output electron beam’s parameters, along with a demonstration of the scheme’s adaptability to create spectra suitable for other critical applications.

3.2.1 Reference model and testing requirements

Before performing any simulations, it was important to select an appropriate radiation model of the Outer Van Allen belt's electron population to be used as a reference for our study. Deemed the most suitable due its accessibility and comprehensive coverage was NASA's AE8-MAX, see Figure 3.8, which is a global empirical model based on data sets from 24 satellites that extend over approximately two solar cycles and provide averages of energetic electron fluxes for solar maximum [153].

This model has been successfully utilised for decades, however it is important to acknowledge that it does possess some inherent limitations. Namely, it provides the long-term average environment for only one solar activity condition and it does not cover populations below 0.1 MeV, meaning the particle flux for this portion of the spectra has to be interpolated and as result, is likely overestimated [154].

Another key task to perform in advance of the study was to establish any pre-existing parametric requirements for electron beams to be used for space radiation hardness assurance testing of electrical components. These would then be employed alongside the energy spectra displayed in Figure 3.8 to set out a criteria for the simulated LPA generated electrons to try and meet.

Upon investigation, industry standards such as NASA MIL-STD 883 (1019.4 for total ionizing dose (TID) testing and 1020.1-1023.2 for dose rate testing) [155] and ESA ESCC [156] recommend a total fluence uncertainty of $\approx 10\%$, which is directly impacted by the total charge and energy spectrum stability of the accelerated electron beam. There is also a guideline listed for beam uniformity [155], stating that it must be maintained to within $\approx 10\%$ in order to ensure consistent radiation exposure. It is important to note that because we are working with broadband energy beams, this level of uniformity has to be present at all energy levels.

Although not explicitly stated within the aforementioned requirements, the beam's divergence is another important quantity to consider as trapped radiation belt electrons can be both strongly divergent and quite directed, for instance during solar energetic particle events [157]. Linac beams, currently utilised for radiation testing, are unable to replicate this diverse beam divergence profile and instead deliver a more focused beam to components ($\theta < 10^\circ$) which ensures acceptable uniform radiation exposure across the test sample. Thus for comparability, we aimed to produce beams with divergence below this threshold.

For other quantities such as the dose rate, there appears to be no general requirements and the acceptable error range is instead defined on a case by case basis depending on the type of effects being studied (TID, single event effects etc.) and the target device. More generally, it is essential to consider the operating conditions of any device being tested, such as bias or temperature, as this can significantly impact the failure rate of electronic systems [158, 159].

Ultimately, it is important to state that although the aforementioned specifications were used for the following study, no specific testing standards exist for LPA generated beams at present due to their infancy within the field of research. Therefore, until such standards have been developed, we chose to look at these existing standards for comparison.

3.2.2 Identification of acceleration regime

The kHz laser parameters in the simulation study were selected to match those achievable by the Coherent Legend Elite Duo He+ laser located within the Strathclyde kHz lab. This made the study more broadly applicable as the Legend Elite Duo is an industrial grade system that is widely available and possesses parameters analogous to many modern day kHz lasers. It also paved the way for any meaningful results that emerged from the simulations to be experimentally verified via an in-house campaign.

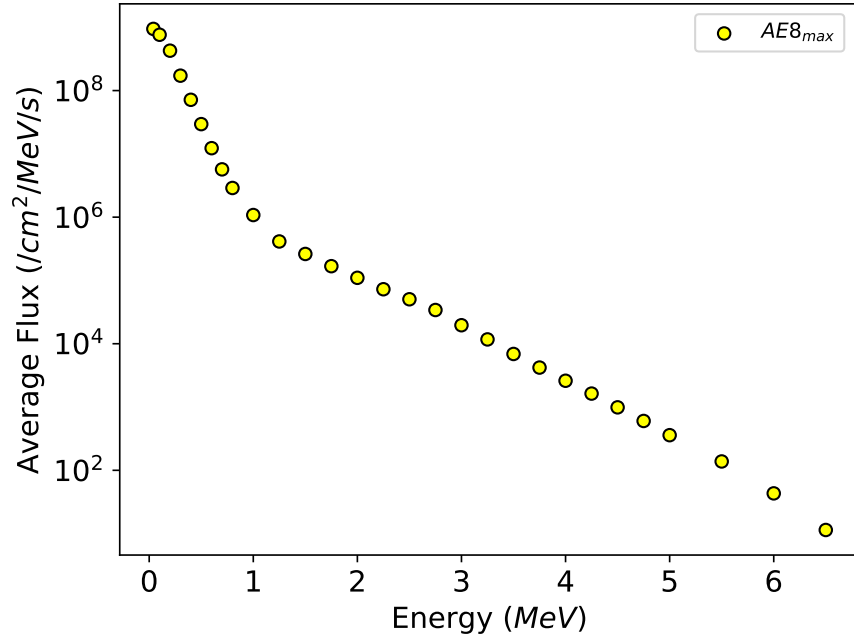


Figure 3.8: AE8-MAX trapped electron model. Empirical model of the average flux of the trapped electron population in the Outer Van Allen belt as a function of energy.

More specifically, an 800 nm Gaussian laser pulse was initialised with an energy of $W_L = 13$ mJ and pulse duration of $\tau = 42$ fs (FWHM). The simulation box size was set at $60 \times 60 \mu\text{m}$, with the cell size in z , $D_z = 5 \times 10^{-8}$ m, the cell size in r , $D_r = 5 \times 10^{-8}$ m and the number of modes used, $N_m = 2$. The number of particles per cell (PPC) along each axis were as follows - $p_{nz} = 2$, $p_{nr} = 2$ & $p_{nt} = 4$. A $150 \mu\text{m}$ symmetrical Gaussian gas profile was generated in the simulation, beginning at $z = 30 \mu\text{m}$ and ending at $z = 180 \mu\text{m}$. The laser pulse was set to propagate for $240 \mu\text{m}$ along z , with the focus set at $z = 90 \mu\text{m}$, which corresponded with the rising edge of the gas jet profile.

Convergence scans were performed as a control test in the PIC simulations, where finer grid sizes, reduced time steps, increased PPC, and a greater number of modes were employed to check that results stayed consistent at these higher resolutions and therefore were in fact physical and not the result of numerical artifacts. In this instance, the PPC selection was made as it allowed for the

electron beam dynamics post-ionisation to be adequately modelled such that accurate energy spectrum, divergence and charge values were produced at the minimum computational cost. Similarly, 2 azimuthal modes were sufficient to resolve the amount of asymmetry generated by these laser-gas interactions and provide converging results with higher N_m values.

The focal spot size was varied throughout, to deliver different peak intensities on target, in an attempt to achieve the required electron beam properties. The lower limit for the laser waist, w_0 , was set at $3\ \mu\text{m}$ as although the theoretical limit is defined by the diffraction limit [160], a number of practical factors such as lens imperfections and the beam quality make achieving a laser waist below this value extremely challenging.

Similarly, both the gas density and the type of gas utilised in the underdense laser-plasma interaction were altered. An upper limit for the neutral gas density was set at $n = 9 \times 10^{19}\text{ cm}^{-3}$. Although this value is more accessible for higher atomic number, Z , gases and could well be exceeded in an experiment with a specialised vacuum pumping system installed, it was desirable to try and reproduce the space radiation spectra at lower gas density values that are more easily achieved and manageable to work with in a practical setting. This same principle applies for the focal spot size, where higher values were deemed as preferable.

The parameter range covered across the study is as follows: the focal spot size was varied between $3\ \mu\text{m}$ and $9\ \mu\text{m}$ in increments of $2\ \mu\text{m}$ whilst keeping the gas density constant. Then, the gas density was varied between $1 \times 10^{19}\text{ cm}^{-3}$ and $9 \times 10^{19}\text{ cm}^{-3}$ in increments of $2 \times 10^{19}\text{ cm}^{-3}$ whilst keeping the focal spot size fixed. This was repeated for each gaseous substance tested in the simulations.

Initial simulations employed hydrogen (H) as the gas of choice, based on its ability to provide superior performance for kHz LWFA with few-cycle, low energy laser pulses through mitigating spatiotemporal distortions imprinted on the laser

field via ionisation-induced defocusing (see Section 2.3.2) and temporal blue-shifting (see Section 2.4.2) [161]. However, no significant electron acceleration was able to be achieved, even when simulating at our defined focal spot and gas density limits. This is because at the energy and pulse duration values we operated with, it was not possible to reach the relativistic self focusing threshold ($P/P_c \approx 1$) in Hydrogen, which as touched upon in 3.1.3, is a requirement for hundreds of keV/MeV scale electron acceleration in any regime for current kHz lasers.

In order to compensate for this, one could increase the density of Hydrogen gas, however this is experimentally challenging in lighter gases due to slower pumping speeds [162] and the higher backing pressures required [161], or move towards higher Z gases. This allows for the release of significantly more electrons upon ionization which in turn, increases both the electron density, n_e , and plasma frequency, ω_p , and reduces the critical power for self focusing, P_c , as shown in Equations 2.22 & 2.42.

Helium (He) was then trialled in the simulations, leading to double the plasma density on-axis, when compared with hydrogen, after the liberation of its two electrons via ionisation. This interaction was considerably more successful at electron acceleration, with the optimized space radiation spectra for this gas displayed in Figure 3.9. This was generated by focusing the laser to a spot size of $5 \mu\text{m}$ whilst operating at a gas density of $9 \times 10^{19} \text{cm}^{-3}$. This caused the laser to self-focus to an $a_0 \approx 1.7$, leading to electrons being self-injected into several bubble shaped cavities that form behind the driving laser pulse. These electrons are then accelerated over $\approx 150 \mu\text{m}$ to generate a beam with large momentum spread. When compared with the AE8-MAX shown in Figure 3.8, the spectra is near identical apart from the very upper end where electrons with energy greater than 7 MeV, that are not present in the trapped electron model, are produced.

However despite the promising spectra, the proposed setup does have two

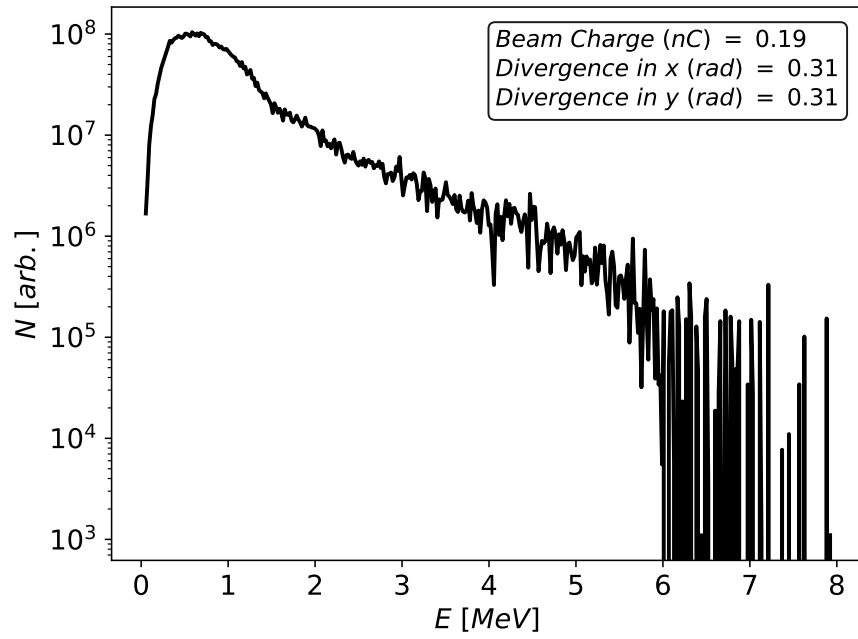


Figure 3.9: Optimised electron energy spectra - Helium. Accelerated electron energy spectra generated by focusing our nominal laser pulse to a spot size of $5 \mu\text{m}$ whilst operating at a gas density of $9 \times 10^{19} \text{ cm}^{-3}$ in Helium.

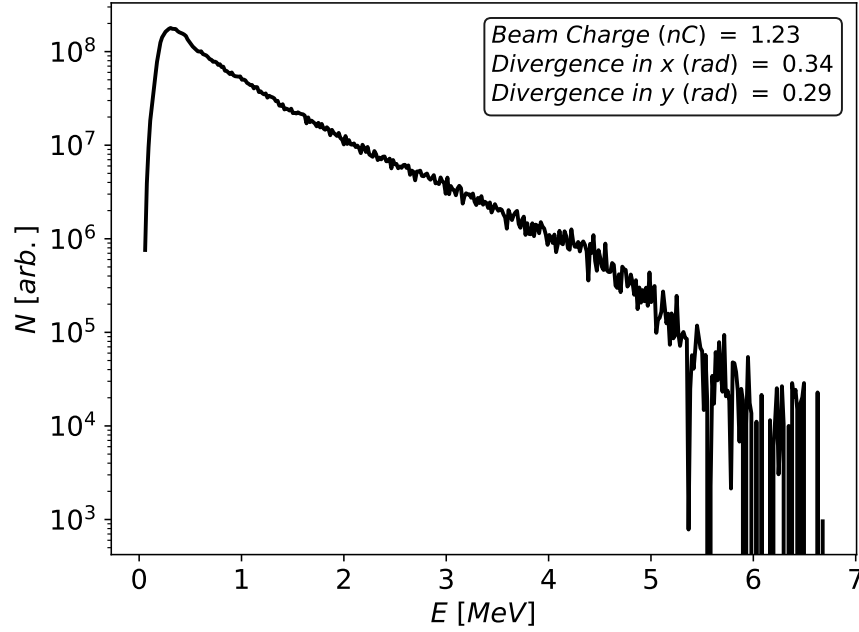


Figure 3.10: Optimised electron energy spectra - Argon. Accelerated electron energy spectra generated by focusing our nominal laser pulse to a spot size of $5 \mu\text{m}$ whilst operating at a gas density of $3 \times 10^{19} \text{ cm}^{-3}$ in Argon.

clear areas for improvement. Firstly, it was only accessible at our defined upper gas density limit and as mentioned previously, it is beneficial to be able to recreate the requisite electron beam at lower densities for experimental accessibility. Secondly, the accelerated e-beam only has 0.18 nC worth of charge which is three magnitudes smaller than the 100 nC beams laser-solid interactions can give rise to [163], admittedly using higher power lasers. Despite this very effectively reducing the peak flux level, it is in fact advantageous to increase this quantity by roughly an order of magnitude [164], assuming kHz repetition rate, as this would enable the flexibility to test against each type of major radiation based effect, both independent and cumulative, and allow for say the daily electron fluence impinging on satellite electronics situated in the radiation belt to be reproduced in less than an hour, further increasing the usability and efficiency of this testing method.

To address both of these factors, Argon (Ar) gas was exploited as each argon atom releases eight electrons assuming immediate ionization of argon to Ar 8+. This meant a far lower gas density was required to produce the same plasma density on-axis as in the Helium case. Furthermore, the larger quantity of plasma electrons present were able to be successfully converted into higher accelerated e-beam charge as exhibited in Figure 3.10, which showcases the optimized space radiation spectra for this gas. This was generated by focusing the laser to a spot size of $5\text{ }\mu\text{m}$ whilst operating at a gas density of $3 \times 10^{19}\text{ cm}^{-3}$. We can see that this interaction succeeds at increasing the charge by a factor of ≈ 10 to 1.23 nC, as well as creating a spectra with no electrons $> 7\text{ MeV}$ and where electrons with higher energies ($> 4\text{ MeV}$) are generated more consistently when compared with the spikes visible in the Helium case.

3.2.3 Characterisation of acceleration regime

Having now identified a suitable accelerator configuration, it was important to properly characterise the laser-plasma interaction taking place in this parameter regime and discern the acceleration mechanisms being exploited. This is of particular importance for the practical implementation of the scheme as certain methods of acceleration are more/less stable than others, which will directly impact on the overall stability of the accelerator.

To examine the laser's evolution over the course of the simulation, snapshots of the pulse's normalised vector potential, a_0 , waist and duration were recorded every $2.5\text{ }\mu\text{m}$ along the propagation and then plotted in Figure 3.11. From initialisation till $\approx 5\text{-}10\text{ }\mu\text{m}$ before the vacuum focal position ($z = 90\text{ }\mu\text{m}$), the pulse behaves as expected, with the waist gradually decreasing towards $5\text{ }\mu\text{m}$, the intensity steadily increasing towards an a_0 value of ≈ 0.591 and the duration remaining relatively constant. However as it approaches the nominal vacuum focus position, the number of plasma electrons on-axis starts to rapidly increase, with the ion-

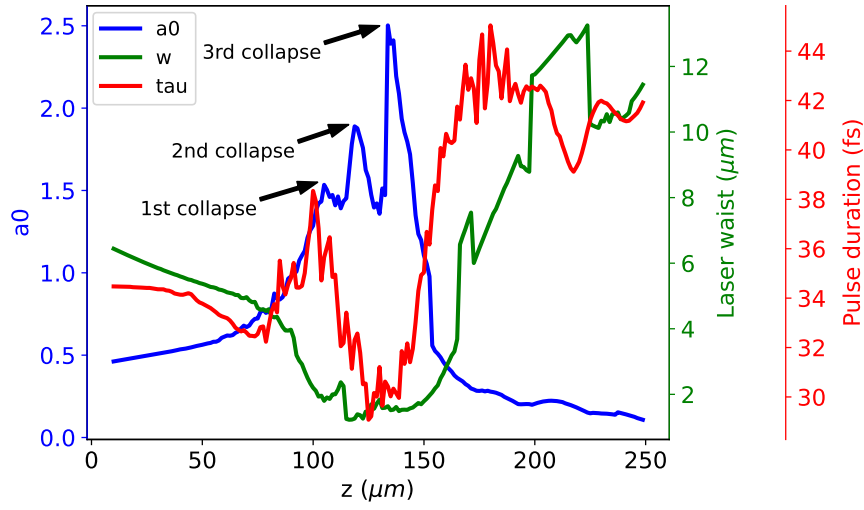


Figure 3.11: Laser pulse parameter evolution along the propagation. The evolution of the laser’s normalised vector potential, a_0 (blue), waist, w (green), and duration, τ (red), as the pulse propagates along the z direction for our optimised Argon case. Note the laser pulse collapse points at $z = 105, 120$ & $130 \mu\text{m}$.

ization of Argon generating a peak electron density of $n_e = 2.4 \times 10^{20} \text{ cm}^{-3}$. This led to a drastic reduction in the critical power for self focusing (see Equation 2.42), $P_c = 126.29 \text{ GW}$, which our laser power, $P_L = 296.53 \text{ GW}$, easily exceeds. Subsequently, the pulse can be seen undergoing severe self-focusing towards a spot size of 1λ which ultimately resulted in laser pulse collapse [121] at $\approx z = 105 \mu\text{m}$.

Successive collapses are found after the first one, as seen in Figure 3.11 at $z \approx 120$ & $130 \mu\text{m}$, as the pulse alternates between strong focusing and strong diffraction, depleting its energy to the plasma each time (see Equation 2.57) [165]. After the first collapse, the pulse duration drops sharply from 38 fs to 29 fs. This is an example of the interconnectivity between laser pulse self focusing and self compression, which are analogous processes but for the waist and duration respectively [166]. Whilst we begin at an amplitude at which self-compression is negligible, the effect of self-focusing is able to push the amplitude up to the

point at which appreciable compression begins to occur, and this in turn further influences the self-focusing and creates even more intense optical fields at the point of collapse. Following the third collapse, the laser pulse's energy has been substantially depleted, such that it now sits well below the critical power for self-focusing, and so we see the pulse remnant diffract and decompress away rapidly. It is important to note that these collapses cause the pulse profile to become highly aberrated and therefore accurate evaluations of its width and duration are difficult to make in these regions, hence the many fluctuations present in Figure 3.11.

For further insight, the laser's E_x field at various points along the propagation is shown in Figure 3.12, with the initial Gaussian pulse profile shown in a). In b), the laser's profile after entering the gas jet is displayed, and we can see that the pulse envelope has begun to contract transversally as expected. Subplot c) shows the electric field at the second laser pulse collapse. As well as continuing to contract, the pulse by this point has become significantly distorted by the formation of several plasma wakes, with the frequency components becoming separated along its extent. Finally, subplot d) shows the fields immediately after the final pulse collapse. Here we can see a pulse that has been further modulated due to scattering processes and laser spectral changes begin to expand transversally.

Each laser pulse collapse is accompanied by a burst of ponderomotively ejected electrons with energies spanning over a few MeV. Figure 3.13 displays the directional components of the electron momentum shortly after the second collapse, with the forward distribution shown in a) and the transverse distribution presented in b). For reference, the laser is polarised in the x direction.

We see that the bulk of these electrons are accelerated in the z (forward) direction, and are generated via laser wakefield acceleration and ponderomotive expulsion. This is examined in Figure 3.14 which shows the 2D longitudinal wakefield amplitude and on-axis 1D wakefield just prior to the second collapse.

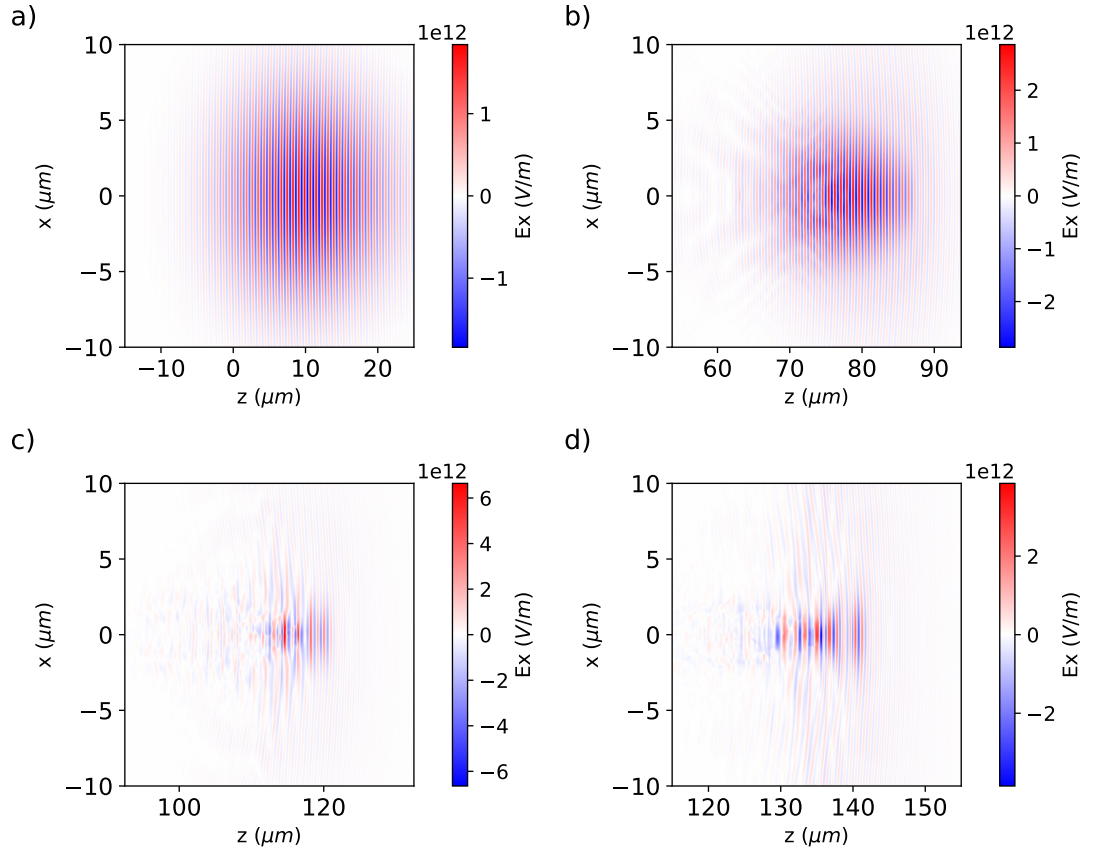


Figure 3.12: Laser pulse E-field spatial and intensity evolution along the propagation. Snapshots of the laser E-field (red/blue) at progressive distances along z . a) $z = 0 \mu\text{m}$, initial Gaussian pulse profile. b) $z = 80 \mu\text{m}$, entering the gas jet. c) $z = 120 \mu\text{m}$, second laser pulse collapse. d) $z = 140 \mu\text{m}$, after the final pulse collapse. The laser pulse is polarised in the x direction.

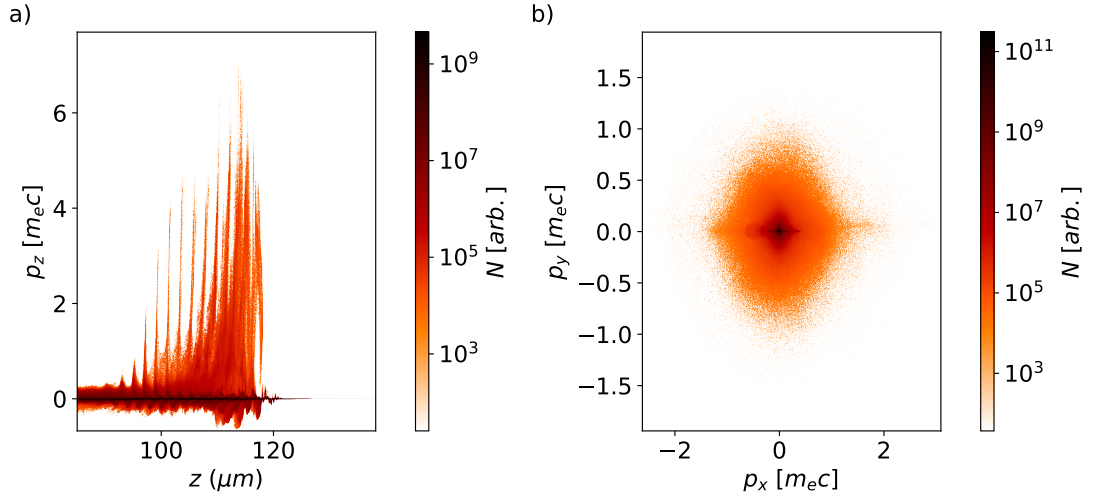


Figure 3.13: Electron momentum after second laser pulse collapse. Directional components of the electron momentum after the second pulse collapse at $z = 120$ μm, where a) shows $z - p_z$ and b) shows $p_x - p_y$.

Here we see evidence of wakefields on the order of TV/m forming along the z axis, ≈ 1.5 μm in length, as a result of the extremely high electron density on-axis (see Equation 2.51). Therefore, although the lifetime of fields of this magnitude is very brief and the dephasing length for this set of laser-plasma parameters results in the acceleration length being drastically reduced, only momentary injection into these wakes is required to rapidly accelerate electrons up to the MeV scale. Another clear observation is that the majority of electrons possess a small amount of transverse momentum. This originates from these particles oscillating in the laser's electric field post ionization.

Taking a more macro view, self-injection in an LWFA has the advantage of being a comparatively straightforward injection method to implement experimentally. However, its reliance in this case on self-focusing and self-compression processes in a high density gas target constitutes a highly non-linear interaction. This can lead to increased fluctuations in the accelerated beam parameters and so it is essential to confirm whether this scheme can still produce an acceptable

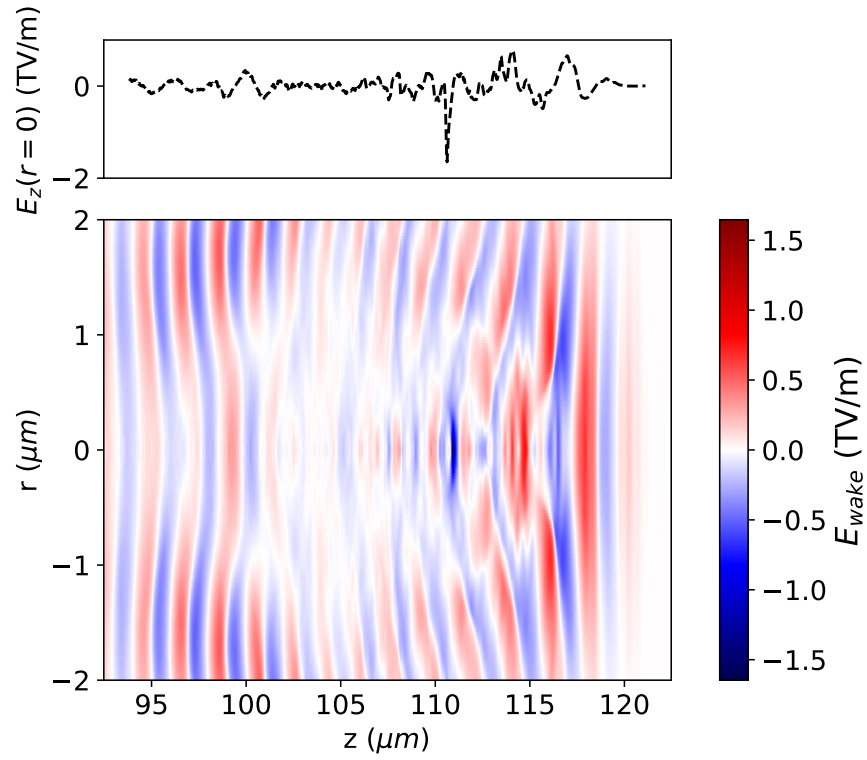


Figure 3.14: 2D longitudinal wakefield and on-axis longitudinal 1D wakefield prior to the second pulse collapse. Illustrative 2D longitudinal wakefield, E_{wake} , and on-axis longitudinal wakefield, E_z , snapshot showing the formation of TV/m scale fields.

output when subjected to parametric variations.

3.2.4 Stability of the accelerator

To establish the overall stability of the proposed accelerator, four key experimental parameters - laser energy, laser focal spot size, laser duration and gas density, were individually subjected to variations in the simulations. The values were as follows: $W_L = 13 \pm 5 \%$ mJ, $w_0 = 5 \pm 10 \%$ μm , $\tau = 35 \pm 10 \%$ fs, $n = 3 \times 10^{19} \pm 10 \%$ cm^{-3} . It is important to state that these simulated fluctuations were selected as worst case scenarios for each respective parameter and are not what would likely be encountered during an experiment. For instance, the Coherent laser's energy stability is listed at 0.22 % rms in its documentation [167].

Upon completion of the simulations, three key output beam parameters - energy spectrum, beam charge and beam divergence, were calculated for each instance and plotted in Figure 3.15. We found that the energy spectrum is well maintained for low and medium energy electrons, however the fluctuations do result in high energy particles (≥ 7 MeV) being produced. These electrons are scarce across all 8 simulated cases suggesting they may in fact be numerical noise, however if physical, they would have a negative impact on delivering the desired total fluence, as described in 3.2.1, and the dose depth profile as they would increase the range of the electron beam such that the shot-to-shot variability will be higher at deeper layers within the electronic device being tested.

The divergence and distribution of the electrons is shown to be maintained within $\pm 10\%$ and is therefore not prone to significant shot-to-shot variability and will have minimal impact on the dose depth profile. However, when examining the beam charge, the maximum shot-to-shot variability observed is 18.7%, which as with the potential energy spectrum inconsistencies, would have an appreciable effect on the total dose and dose depth profile.

Fortunately, since we integrate over multiple accelerated beams in this pro-

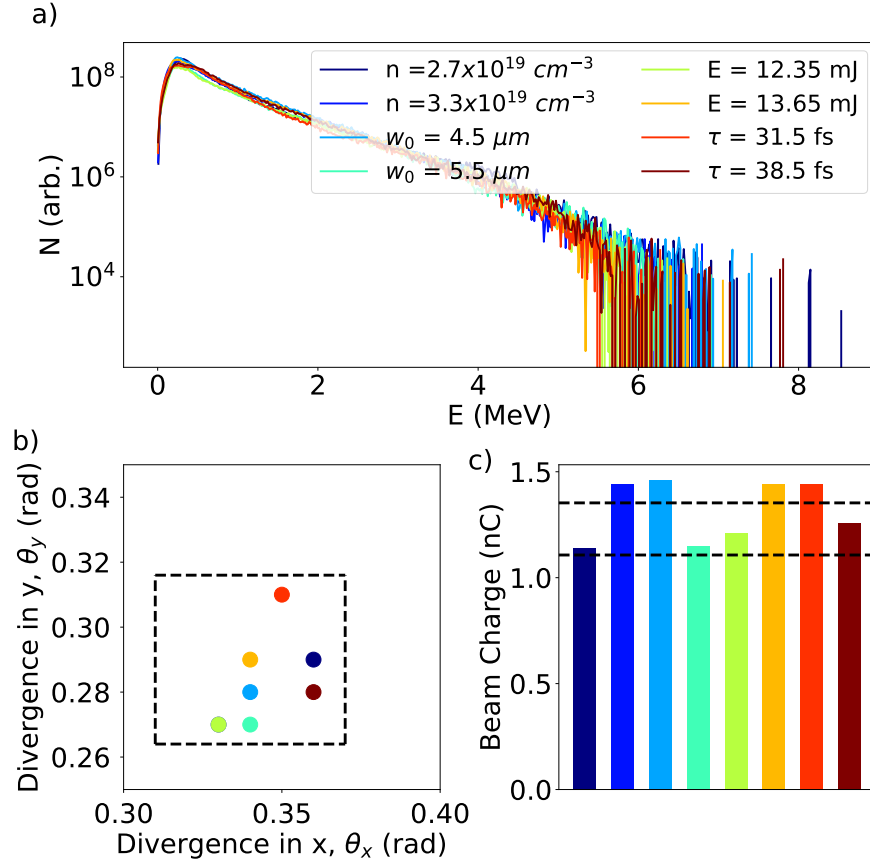


Figure 3.15: Stability of the accelerated electron beam. Three key output electron beam parameters - a) energy spectrum, b) beam divergence and c) beam charge, displayed for several simulated cases (see legend in a) for assigned colours) where key experimental parameters have been varied to allow the overall stability of the accelerator to be diagnosed. The black dashed line in subplots b) and c) signifies a 10 % deviation in parameter value from our mean nominal, non-fluctuating case.

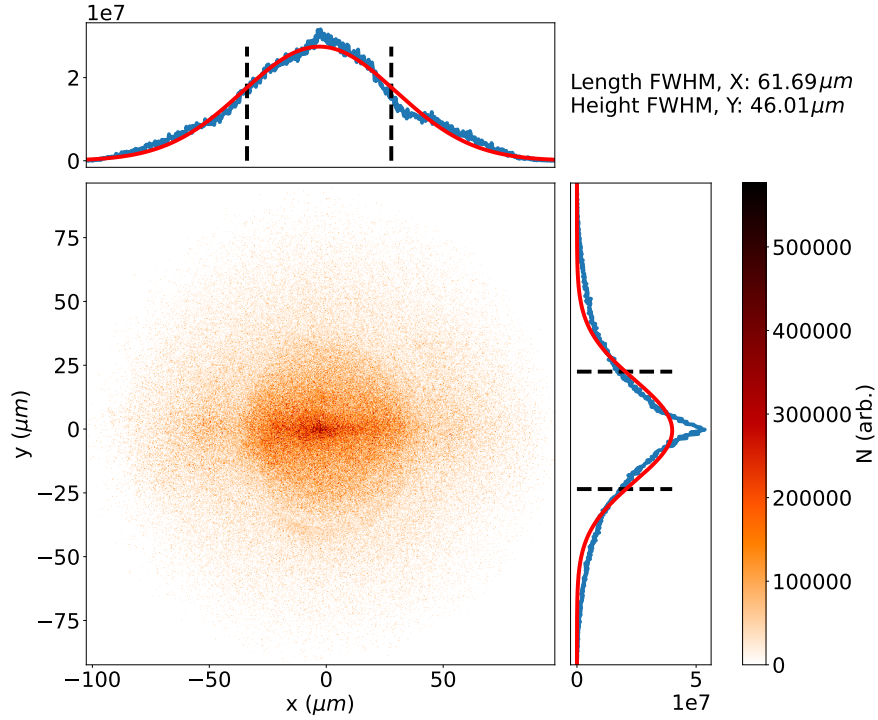


Figure 3.16: Accelerated beam profile - Linearly polarized laser. Transverse real space distribution of the accelerated electron beam for our optimised Argon case parameters utilising a linearly polarized pulse. The horizontal (x) and vertical (y) line-outs are plotted on the top and side (light blue), and the FWHM size is determined by fitting a Gaussian function (red).

posed testing method, the final dose depth profile should average out to the expected dose profile as long as a low dose per pulse is maintained to ensure hundreds or thousands of beams are required to reach the desired dose in the device under testing (DUT). The dose per pulse can be significantly reduced or increased by varying the source-to-target distance without needing to change the laser repetition rate. This averaging out effect successfully mitigates against shot-to-shot fluctuations, even in the case of the charge variation where we possess a large difference between the maximum and minimum value.

3.2.5 Electron beam parametric improvement

As touched upon in 3.2.1, another one of the central requirements for space radiation hardness testing is that the dose variation over the target volume is restricted so that all points within the volume receive the prescribed dose within a tolerance range of $\leq 10\%$.

Two quantities which directly affect the uniformity of the dose in target volume are the symmetry and flatness of the electron beam, with symmetry referring to the similarity of dose distribution on opposite sides of the beam's central axis, whilst flatness describes the uniformity of dose across the beam's field [168]. The transverse real space projection of the accelerated beam is plotted in Figure 3.16 to allow for these two variables to be diagnosed.

Starting with the beam symmetry, we observe an electron bunch that possesses a degree of asymmetry, with a higher portion of particles distributed along the laser polarization direction, x , than the perpendicular plane, y . This is a feature that has previously been observed in laser-plasma interactions and is now well understood [169], arising from electrons generated via optical field ionisation possessing the majority of their initial momentum in the polarisation direction (see Figure 5.13), which in this case is x .

In an attempt to improve this unevenness, a circularly polarized laser pulse was initialised in the simulation, replacing the linearly polarized pulse as our LPA drive laser, and the simulation was re-ran with our nominal parameters - $W_L = 13$ mJ, $\tau = 42$ fs (FWHM), $w_0 = 5$ μm and $n = 3 \times 10^{19} \text{ cm}^{-3}$ in Argon gas. It is important to note that for the same laser intensity, the a_0 value for a circularly polarised pulse is $\sqrt{2}$ times smaller than for a linearly polarised pulse. This move was prompted by [170], who found that circular polarized light allowed them to create a more radial electron beam profile that focused better into a symmetric central spot, rather than the Gaussian/Lorentzian transverse real space profile associated with linear polarized light. The resulting transverse

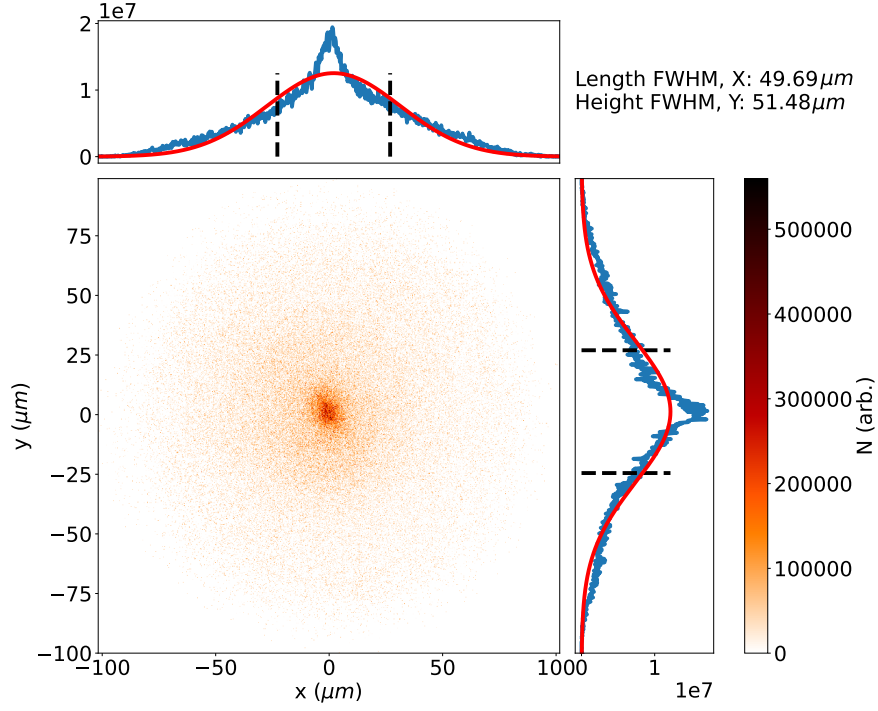


Figure 3.17: Accelerated beam profile - Circularly polarized laser. Transverse real space distribution of the accelerated electron beam for our optimised Argon case parameters utilising a circularly polarized pulse. The horizontal (x) and vertical (y) line-outs are plotted on the top and side (light blue), and the FWHM size is determined by fitting a Gaussian function (red).

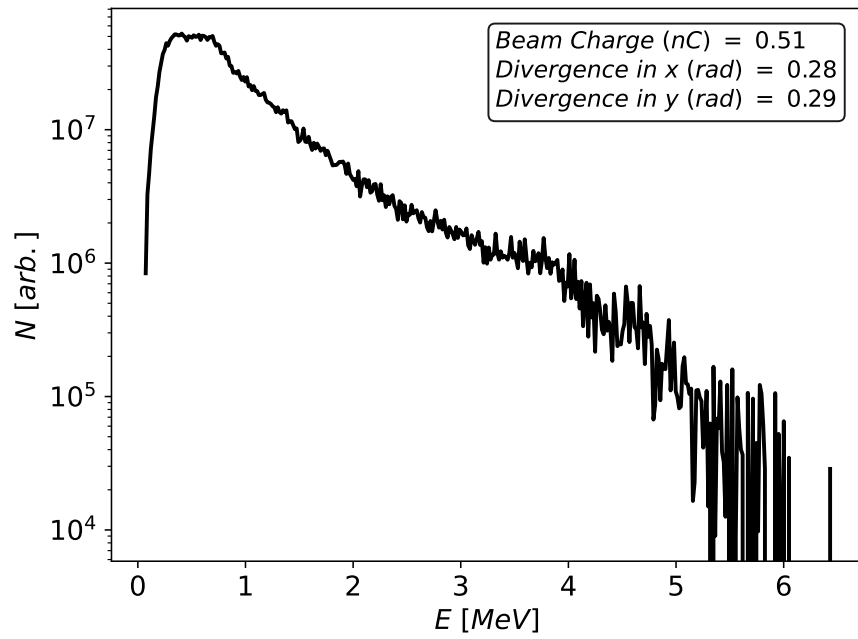


Figure 3.18: Accelerated electron energy spectra - Circularly polarized laser in Argon. Accelerated electron energy spectra generated by focusing a circularly polarized laser pulse with our nominal parameters to a spot size of $5 \mu\text{m}$ whilst operating at a gas density of $3 \times 10^{19} \text{cm}^{-3}$ in Argon.

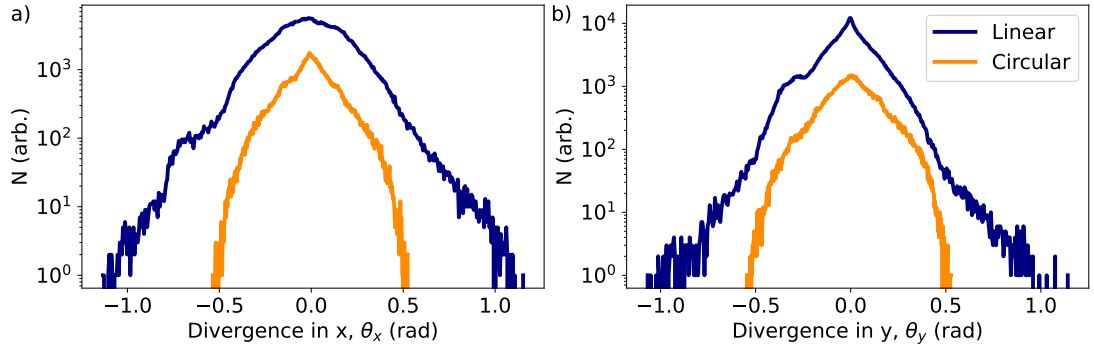


Figure 3.19: Accelerated electron beam divergence comparison. Histogram comparing the divergence of the electron beam generated by a linearly polarized pulse (blue) and a circularly polarized pulse (orange) for our optimised Argon case parameters. The x (a) and y (b) components of the divergence are displayed separately.

real space projection of the accelerated beam is plotted in Figure 3.17, followed by its associated energy spectrum in Figure 3.18.

Clearly visible is the beam's improved symmetry around the central axis when compared with the linear polarization case shown in Figure 3.16, with the disparity in x-y FWHM length reduced from $15.68 \mu\text{m}$ to $1.79 \mu\text{m}$. This enhanced symmetry is then maintained in the beam divergence, displayed in Figure 3.19, where we see a reduction in the spread of x and y particle divergence across the beam, and closer absolute divergence values - $\theta_x = 0.28$ and $\theta_y = 0.29$, when compared with $\theta_x = 0.34$ and $\theta_y = 0.29$ in the linear polarisation case. Also, the electrons now appear to have no clear directional bias. Most importantly, the accelerated electron spectra is still comparable to our optimized space radiation spectra (see Figure 3.10) despite the change in orientation of the oscillating electric fields in our laser beam, with the principal difference being the reduction in charge caused by a lack of low energy electrons ($< 1 \text{ MeV}$).

Looking now at the beam flatness, it is evident that for both the linear and circularly polarized cases, there is a lack of uniformity across the electron bunch's extent, with the profile resembling a low variance Gaussian distribution as oppose

to a flat-top shape. The most simplistic way to create a uniform intensity distribution would be to pass the beam through an aperture which blocks all but the central, and most uniform portion of the beam. However, this would ultimately be problematic as a large fraction of the electrons would be discarded, distorting the energy spectrum, and the resulting beam would still have a notable falloff in intensity from the centre to the edge.

A more nuanced approach would be to create a beam forming system by first passing the beam through a scattering foil, aluminium and polyethylene are popular candidates, which flattens the central part of the Gaussian and scatters the beam to a larger area [171], and then the scattered electron bunch can be sent through a collimator to confine it to the area of the desired irradiation field [172]. Despite decades of electron beam forming systems being utilised for various applications in medicine, science and industry, the scheme still possesses some deficiencies, particularly beam straggling, which is triggered through collisions with the atoms in the target foil, and causes degradation in the energy spectrum. This effect is most pronounced on electrons with energy ≤ 0.2 MeV, however for many key radiation-based effects in spacecraft instruments such as internal charging and total ionizing dose effects, the impact of fluctuations in the spectrum of these low energy electrons is minimal and therefore not a showstopper for testing purposes.

Finally, magnets are now being used in electron beam systems to provide precise control over the beam's trajectory and profile, allowing for refocusing, collimation, and correction of aberrations to improve beam symmetry [173, 174]. This approach has many advantages over the aforementioned methods as it is inherently non-destructive, can provide smooth continuous fields that lower emittance growth and offers active control and adjustment to correct asymmetries dynamically during operation.

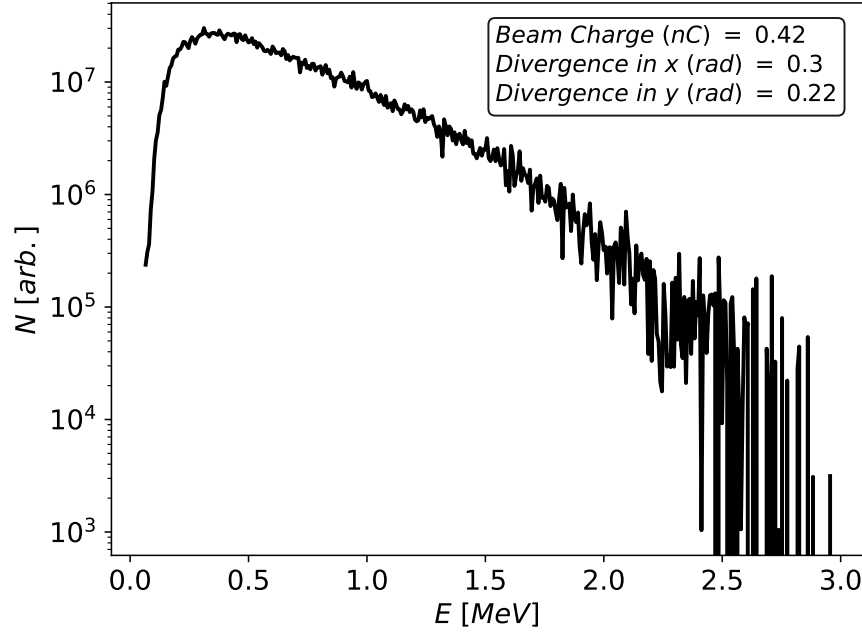


Figure 3.20: Favourable skin cancer therapy electron energy spectra. Accelerated electron energy spectra generated by focusing our nominal laser pulse to a spot size of $9 \mu\text{m}$ whilst operating at a gas density of $9 \times 10^{19} \text{cm}^{-3}$ in Argon.

3.2.6 Accelerator flexibility

Finally, it is important to highlight the flexibility of this proposed accelerator setup in that spectra ideally suited to several other applications can be achieved solely by tweaking the input parameters laid out in subsection 3.2.2. For instance, the treatment of surficial skin cancers could be greatly improved by exploiting broadband electron beams from LPAs with energies up to 3 MeV as they combine low deep tissue damage with a short dose fall-off [175]. Here we are able to generate such a spectra by simply increasing the gas density to $9 \times 10^{19} \text{cm}^{-3}$ and focusing the laser pulse less tightly to a spot size of $9 \mu\text{m}$, as displayed in Figure 3.20.

Furthermore, the quick and local deposition of energy imparted by these beams would be well suited for fast ignition, one of the most promising approaches

to reach a high gain in inertial confinement fusion [176], further underlining the accelerator’s versatility.

3.2.7 Summary

The work summarised in this subchapter constitutes an important step in demonstrating the ability of a contemporary kHz laser to replicate the energy spectra and flux of electrons located within the Outer Van Allen belt by way of a laser-underdense plasma interaction. The simulated parameter scan required to determine an acceptable interaction regime was presented, followed by the delineation of the acceleration mechanism being exploited. The stability of this potential accelerator was outlined alongside methods for improving key output electron beam parameters, allowing this proposed setup to be employed as a viable alternative beam source to satisfy the growing demand for space radiation hardness assurance testing.

As discussed in Section 3.2.6, the energy range and broadband nature of the electron beam spectra created by this accelerator lends itself to many other applications. A similar parametric scan would be advisable to find the idealised spectra for each individual use case, noting that heightened requirements may be placed on electron stability, for instance greater restrictions on dose variability must be applied for skin cancer therapy when compared with space radiation testing as instead of trying to find the statistical mean dose at instrument failure, the aim is to deposit a specified amount of dose.

An experimental campaign to verify the results laid out in this subchapter, which to date has been prohibited by insufficient radiation protection within the SCAPA kHz lab, is the logical next step in terms of future studies. This would allow for the effect that the multitude of parameters not reflected in our PIC simulations have on the output electrons to be quantified. Inversely, the accuracy of our PIC simulations can be determined and the prefile settings adjusted so as

to simulate a more lifelike scenario.

Chapter 4

Hybrid plasma photocathode acceleration

Injecting electrons into an laser wakefield accelerator to form a ‘witness’ beam can be achieved through various processes, ranging from self-injection [62–64] to more elaborate methods involving local modifications of plasma density [177–180], multiple intense laser pulses [181], using higher ionization levels and dopants [71, 182–184], as well as combinations of these methods [185, 186].

A common feature across these schemes is that electron injection is tightly coupled to the wake excitation process. Fluctuations in the drive laser’s characteristics – such as power, polarization [187], or wavefront quality [188] – along with the inherently non-linear and relativistic nature of the laser–plasma interaction [49], significantly influence the evolution of the plasma wakefield. These coupled dynamics, in turn, affect injection rates and charge yield, resulting in shot-to-shot variations in the properties of the injected electron beam.

Once injected, the LWFA-inherent challenges of diffraction of the laser driver and dephasing with the injected electrons are further complications that can change the longitudinal and transverse phase space of the injected electron beam throughout the subsequent acceleration process [49]. While meticulous control

and fine-tuning of the laser-plasma interaction can improve stability, state-of-the-art systems still exhibit fluctuations in charge, energy, and energy spread on the order of a few to 10% shot-to-shot [186, 189, 190].

However, engineering optimization of LWFA is not the only path towards improved stability. In this chapter, intrinsic compensation mechanisms when using fluctuating LWFA output to power a subsequent electron beam-driven plasma wakefield acceleration (PWFA) [65, 66] refinement stage [75–78, 191–193] are explored. In PWFA, strong plasma waves are excited by the unipolar Coulomb force of an electron beam, but with peak driver fields orders of magnitude lower than in LWFA. This environment enables controlled electron injection via a plasma photocathode [93, 116, 194], where a low-power laser pulse with $I_L \sim 10^{15} - 10^{16} \text{ W/cm}^2$, just above the tunneling ionization threshold, releases electrons directly into the wake, with negligible momentum imparted by the injection laser [169].

The plasma photocathode can thus generate well-defined electron bunches with potentially ultralow normalized emittance $\epsilon_n < 100 \text{ nm rad root mean square (r.m.s.)}$ [93, 116, 169, 194, 195], while the release process and injection rates are decoupled from the driver and wakefield configuration. While staging of plasma wakefield accelerators is usually associated with challenges in capture efficiency and beam quality preservation [196], this work demonstrates that staged electron beam utilization from LWFA into plasma photocathode-based PWFA can yield electron output bunches with significantly higher stability, in addition to substantially enhanced bunch quality. These improvements come at the cost of total electron beam charge, but for applications such as light sources where electron beam quality thresholds are crucial, this is a valuable trade-off [197].

To explore the interplay of key mechanisms in this stabilisation process, high-fidelity, quasi-3D particle-in-cell simulations of the PWFA refinement stage equipped with a plasma photocathode were performed using FBPIC [83]. The electron driver beams utilized in the study encompass a wide parameter space

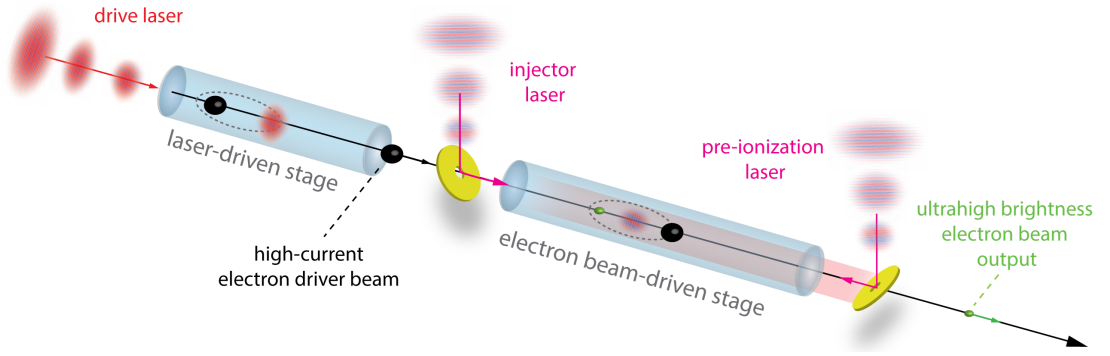


Figure 4.1: Collinear staged plasma photocathode wakefield accelerator in the laboratory frame. An intense laser pulse first generates a nonlinear plasma wave in the initial stage. Electrons injected and accelerated in this region produce a dense but low-quality electron beam (black ball). This beam then passes through a holed mirror (golden ring) and is transferred into a second, pre-ionised stage, where it drives a plasma wakefield accelerator. Within this stage, a plasma photocathode injects a high-quality witness beam, which eventually can be exploited for subsequent applications such as light sources

accessible by contemporary LWFA systems via a variety of methods, including, but not limited to, self-truncated ionization injection [71, 183, 184] or shock-front injection [72, 80, 179, 198, 199].

4.1 Simulation parameters and setup

For the showcase PWFA configuration, the plasma density is set to $n_e \approx 1.12 \times 10^{17} \text{ cm}^{-3}$, corresponding to a cold plasma wavelength $\lambda_p = 2\pi c/\omega_p \approx 100 \mu\text{m}$. The baseline electron driver beam is bi-Gaussian with charge $Q_D = 500 \text{ pC}$, length $\sigma_z = 6.0 \mu\text{m}$ (r.m.s.), radius $\sigma_r = 4.0 \mu\text{m}$ (r.m.s.), mean kinetic energy $W_D = (\gamma_D - 1)m_e c^2 = 500 \text{ MeV}$, where γ_D is the beam's Lorentz factor, relative energy spread $\Delta W_D/W_D = 10\%$ (r.m.s.), normalised emittance $\epsilon_n = 2.0 \mu\text{m rad}$ (r.m.s.), and peak current $I_p \approx 10 \text{ kA}$.

These parameters were not chosen to simulate a specific experiment, but rather to represent an LWFA-like beam with some assumptions made around

the ability to generate a higher charge bunch at the expense of beam quality. A charge value of $Q_D = 500$ pC was chosen to allow the blowout to be efficiently driven at a lower plasma wavelength, where as the emittance value was selected because well executed LWFA shockfront injection typically produces beams with $\epsilon_n = 1.0 \mu\text{m rad}$ (r.m.s.) [94], and so double this value seemed reasonable for a higher charge beam. By choosing these values, neither the longitudinal PWFA resonance condition [97] $\sigma_{z,m} = \sqrt{2}/k_p \approx 22.5 \mu\text{m}$ nor the transverse PWFA matching condition [200] $\sigma_{r,m} = \sqrt[4]{2/\gamma_D} \sqrt{\epsilon_n/k_p} \approx 1 \mu\text{m}$ is fulfilled. This is intentional, to account for non-ideal conditions that may arise from unoptimized beam transport between stages – further complicated by the inherent volatility of the LWFA electron beam output – and to underscore the robustness of the plasma photocathode-enhanced PWFA stage. This choice is also appropriate because, in present day hybrid LWFA→PWFA [201] setups, the electron beams emerging from the LWFA stage are typically intrinsically short, carry multi-kA currents [73, 74, 202], and expand rapidly in the transverse dimension between stages.

Nevertheless, the driver beam’s peak density, n_b , which is expressed as $n_b = Q_D / [e(2\pi)^{3/2}\sigma_r^2\sigma_z] \approx 2.1 \times 10^{18} \text{ cm}^{-3}$, exceeds n_e by more than an order of magnitude. The driver is therefore sufficiently overdense to enter the nonlinear blowout regime of PWFA [203], even under substantial driver charge variations. The blowout regime is characterized by longitudinally uniform linear focusing fields and a radially uniform accelerating field with a linear slope around the wake centre (see Section 2.5.1). This linear field configuration is a prerequisite for preserving the normalized emittance during acceleration [203].

Utilizing intrinsically short driver beams from LWFA, confined to the front of the blowout, is furthermore beneficial for plasma photocathodes, as this eliminates any influence of the driver fields on the plasma photocathode laser field ionization rate, as well as transverse momentum and witness beam emittance growth, over a wide range of possible release positions [116, 169, 204]. The radial electric field

of the driver beam is given by $E_r(r) = (Q_D / [(2\pi)^{3/2} \sigma_z \epsilon_0 r]) \times (1 - \exp(-r^2/2\sigma_r^2))$, and peaks at approximately 67 GV/m. This is sufficient to field-ionize helium from its ground state (low ionization threshold, LIT: 24.6 eV), while leaving its second, higher ionization threshold level (HIT: 54.4 eV) unperturbed. In this study, preionization of the first helium level using an auxiliary laser pulse is assumed, as recently demonstrated experimentally [76–78, 80]. This controlled preionization supports a more stable and robust wakefield formation [76, 77].

This hybrid LWFA→PWFA configuration is then equipped with a collinear plasma-photocathode laser, delayed by approximately 133 fs with respect to the driver centroid. The laser operates with an energy of $W_L = 0.1$ mJ and a pulse duration of $\tau = 30$ fs full width at half maximum (FWHM). It is focused to an r.m.s. spot size of $w_0 = 5 \mu\text{m}$ at the blowout centre – a unique feature of the plasma photocathode that gives rise to several inherent stability mechanisms discussed in this chapter. This setup enables witness bunch generation by controlled ionization of the HIT (He^+) ions.

A diagram of the proposed experimental layout is displayed in Figure 4.1. It is important to acknowledge that inter-stage beam transport and matching, as well as emittance preserving witness extraction from the plasma and subsequent separation from the driver remain challenges for the experimental implementation of this concept, and therefore further advancements are required to extend the scheme beyond a fundamental PWFA testing environment towards a reliable, compact source of ultra-high brightness beams.

Figure 4.2 illustrates the showcase PWFA simulation configuration, showing the plasma photocathode process as the injector laser traverses its focus. Consecutive PIC simulation snapshots are shown in the co-moving coordinate system $\zeta = z - ct_D$, given that the wake propagates at approximately c due to the relativistic velocity of the driver. Here, t_D denotes the driver’s propagation time in the plasma.

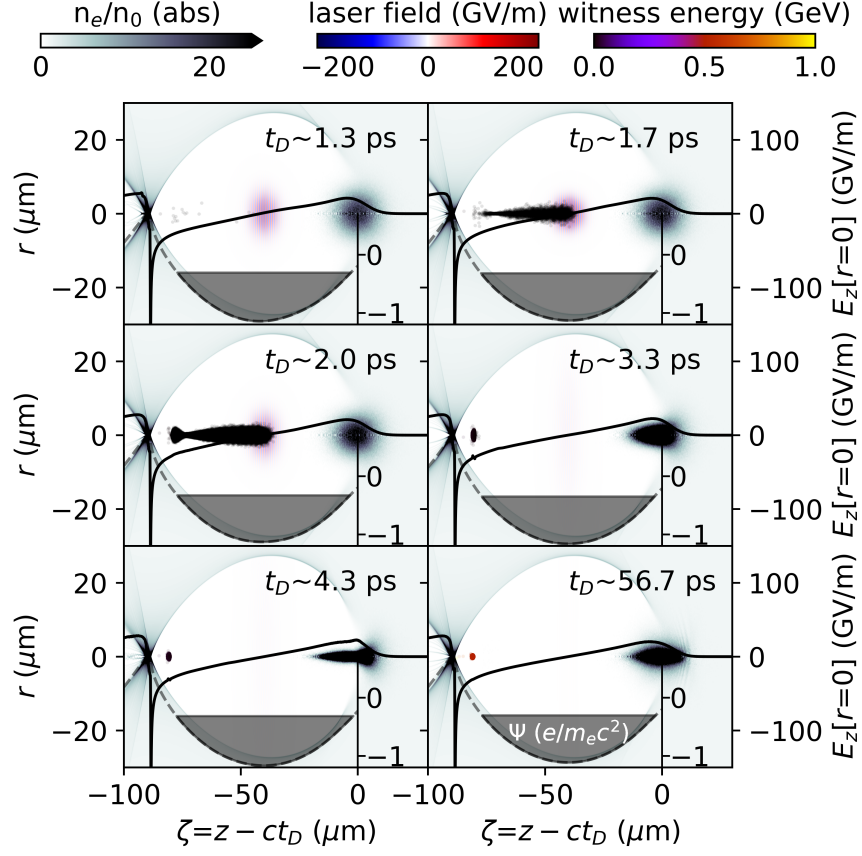


Figure 4.2: PWFA showcase simulation. Simulation snapshots illustrate plasma photocathode injection within a helium plasma wave. The electrostatic pseudo-potential Ψ is shown as a grey dashed line, and the shaded region indicates the longitudinal trapping window ($\Psi - \Psi_{\max} \leq -m_e c^2/e$) within the blowout structure. The corresponding on-axis electric field E_z is plotted as a solid black line. An electron beam drives the wakefield as it propagates to the right. A co-propagating laser pulse (blue to red) ionizes He^{2+} , releasing witness electrons that are color-coded by energy (black to yellow).

The electrostatic pseudo-potential Ψ and the corresponding longitudinal component of the on-axis electric field ($E_z = \partial_\zeta \Psi$) are shown in Fig. 4.2, with the zero crossing located at approximately $\zeta \approx -40 \mu\text{m}$. Electrons are released at rest by the plasma photocathode laser, with the laser delay chosen such that the ionization front aligns with the minimum of the potential well (see Fig. 4.2). As the electrons are born and slip backward in the co-moving frame, the electrostatic pseudo-potential is converted into kinetic energy, accelerating them. If the potential is sufficiently large, it will facilitate trapping of electrons in the plasma wave. 100% trapping efficiency is achieved in the showcase PWFA simulation configuration due to the aforementioned witness electron beam release position and momentum, as well as the driver beam possessing sufficient charge (see Section 4.2).

The trapping condition can be derived from the Hamiltonian of an electron initially at rest ($v_i \approx 0$) in a blowout propagating with phase velocity $v_\phi \approx c$ [205]. In contrast to the convention used in Ref [205], where the pseudo-potential is maximized in the blowout centre, a convention where Ψ has a minimum at the zero crossing similar to [89] was adopted in this study. Accordingly, an electron liberated at co-moving coordinate ζ_{inj} , where the pseudo-potential is Ψ_{inj} , will become trapped in the wakefield if there exists a co-moving coordinate ζ_{trap} such that the condition $\Psi_{\text{inj}} - \Psi_{\text{trap}} = -m_e c^2/e$ is satisfied.

The range of injection coordinates that lead to trapping in the wakefield is given by the condition $\Psi_{\text{inj}} - \Psi_{\text{max}} \leq -m_e c^2/e$, indicated by the gray-shaded region in Fig. 4.2. This relationship between injection position and Ψ independently determines the final co-moving trapping location within the plasma wave for cold released electrons, thus fixing the electric field by which the electrons are subsequently phase-constantly accelerated.

The parabolic symmetry, together with the localized injection dynamics governed by the electrostatic potential and the decoupled nature of the plasma photo-

cathode, form key foundations of the stability mechanisms discussed throughout this chapter. The primary fluctuations and their contributions to the stability process investigated in the following sections are the electron driver energy, energy spread, and current, as well as the plasma photocathode injector laser energy and its spatiotemporal synchronization with the driver. It is important to note that the driver charge density, which is the key parameter responsible for the wakefield excitation, could also be changed by varying the driver length or its transverse extent, but such additional studies would be largely redundant with changing the driver charge.

4.2 Witness charge is fully decoupled from driver beam charge

To reflect shot-to-shot jitter in this setup, simulations were performed with significant variations in the driver beam density by varying its charge Q_D . In contrast, the released and trapped witness charge remains practically completely independent of drive beam charge fluctuations. This insensitivity arises because the witness charge yield is determined by the plasma photocathode injector laser field strength and the helium density, which together set the number of electrons that can be released per unit volume from He^+ ions.

With a baseline dimensionless light amplitude $a_0 \approx 0.0611$, the pulse will tunnel ionize the He^+ ions and will release He^{2+} electrons with $Q_W \approx 4.9$ pC of charge. The simulation results shown in Fig. 4.3(a) reveal that a witness charge capture efficiency of 100% is achieved, as long as the driver charge amounts to $Q_D \gtrsim 290$ pC, corresponding to a current of ~ 5.8 kA, the threshold where the electrostatic potential becomes large enough to trap all electrons released around the blowout centre.

In principle, the wakefields themselves may also contribute to the charge yield

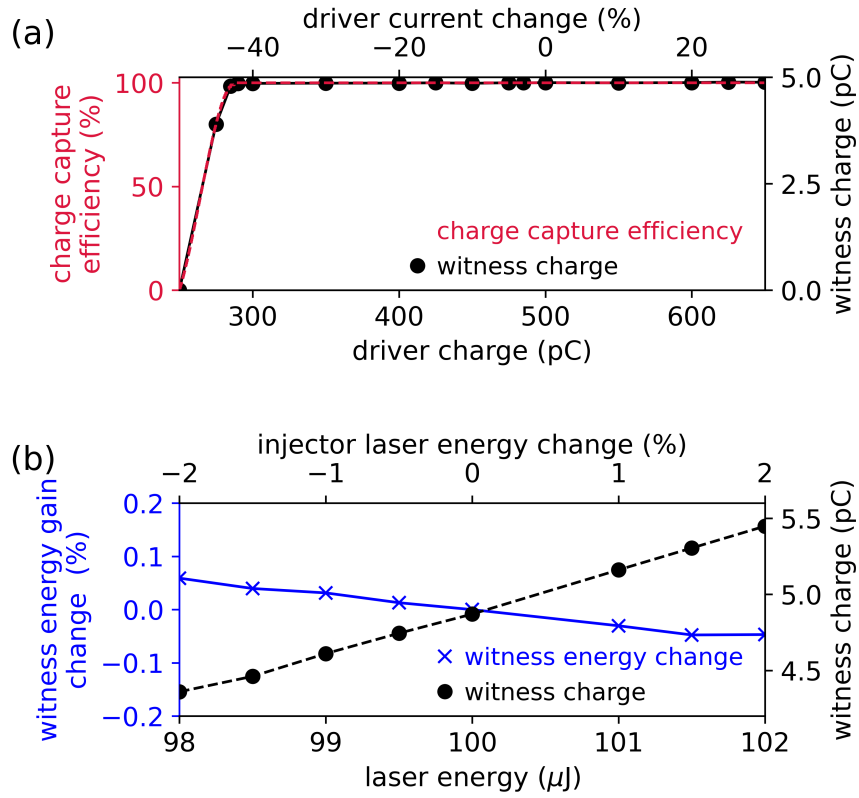


Figure 4.3: Witness charge when driver charge or laser energy varies. The stability of the injected witness charge versus varying driver charge and current is shown in panel (a) at constant laser energy, while panel (b) shows witness energy stability versus injector laser energy variation at constant, baseline driver beam charge.

through superposition with the laser pulse field, leading to a local increase in the electric field amplitude. However, releasing at the blowout centre has the concurrent advantage that not only is the electrostatic potential minimized, but the wakefields are nearly zero in this region [204]. Therefore, the wakefield has a negligible impact on the charge yield. Furthermore, since this showcase uses the second ionization level of helium as the HIT level – in contrast to the first ionization level of He used in, e.g., Ref [116] – the laser intensity is chosen to be significantly higher. As a result, the contribution from shot-to-shot wakefield variation becomes even more marginal. In any case, the produced witness bunch charge is fully decoupled from the driver beam charge over a very wide range.

While charge capture efficiency and stability are natural outcomes of the laser-dominated collinear hybrid LWFA $\rightarrow$ PWFA plasma photocathode setup, these factors present significant challenges in other LWFA or PWFA injection and staging schemes. For example, in plasma density downramp injection in LWFA [72] or PWFA [78, 117, 198, 206], the injection rates are strongly dependent on the strength of the driver, making them highly sensitive to fluctuations of the driver beam from shot-to-shot.

Similarly, during LWFA staging [196] or external injection [207], full charge capture is very challenging, arising from a high sensitivity to alignment between the trajectory of an already relativistic electron beam and the subsequent wake [208]. As a result, spatiotemporal fluctuations and/or angular offsets of the injected electron beam with respect to the wake have far-reaching consequences for the subsequent acceleration dynamics.

While driver beam electron density is the key parameter for wakefield excitation, the corresponding critical parameter for the plasma photocathode injector is the laser intensity. It is therefore worthwhile to investigate the witness beam charge when the laser energy is varied, but the driver beam charge density is assumed constant.

When the injector laser energy is varied by $\pm 2\%$ around this baseline, the corresponding tunneling ionization rates produce witness bunches with charge variation of 4.9 ± 0.5 pC ($\pm 10\%$) as shown in Fig. 4.3(b), right y -axis. The slightly different charge has a very small effect on the energy gain, which due to beam-loading is smaller for higher laser energy than for lower laser energy (see Fig. 4.3(b), left y -axis).

The plasma photocathode laser pulse can either be provided as an intrinsically synchronized pick-off from the LWFA laser system, or delivered via an external, synchronized laser arm. Since a laser energy variation of $\pm 2\%$ is significantly worse than the operational standard achieved by commercial high-power laser systems shot-to-shot, this parameter should be regarded more as a tunable knob for adjusting the witness beam charge – and consequently compound quality metrics such as the witness bunch brightness [209] – rather than a source of jitter. Moreover, various techniques may be explored to normalize the pulse energy at the sub-mJ level by sacrificing input pulse energy, such as cross-polarized wave normalization, reverse saturable absorbers [210], non-linear optical limiters [211] etc.

4.3 Witness energy gain is independent of driver beam energy

Next, the impact of energy fluctuations of the LWFA-generated driver beam was investigated, by varying its central energy from 125 MeV to multi-GeV levels, whilst keeping its current, transverse properties and the injection laser properties fixed.

Figure 4.4 (a) reveals that the witness energy gain rate is linear, and nearly constant across the driver parameter sweep, despite the order of magnitude variation in driver energy. Fundamentally, this is because at sufficiently high

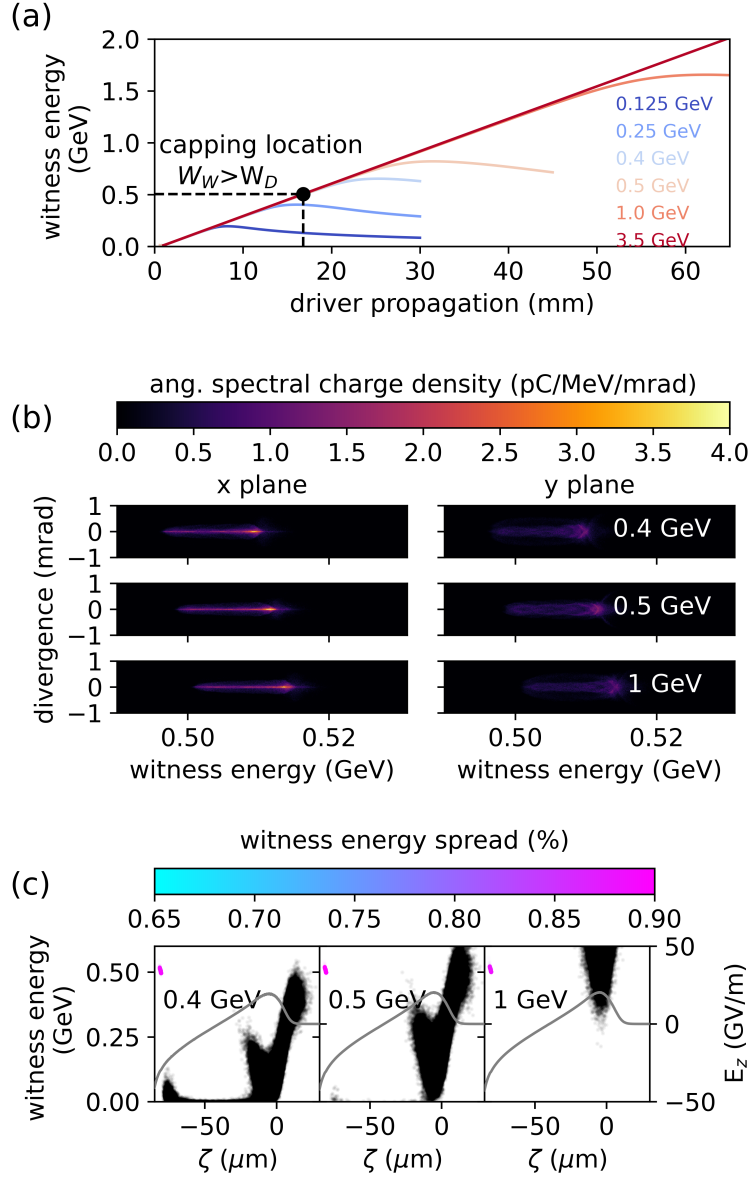


Figure 4.4: Impact of driver beam electron energy (0.125–3.5 GeV) on witness energy gain. Panel (a) shows the energy gain of the witness beam as the driver beam energy varies. Panel (b) presents the spectral charge density projections of the witness beam in both the x and y planes for driver energies of 0.4 GeV, 0.5 GeV, and 1 GeV, at a potential cut-off location indicated by the dashed line in panel (a). Panel (c) displays the energy phase space plots of both the driver (black) and witness beams (left y -axis), color-coded by relative energy spread, corresponding to the conditions in panel (b) and the values contained within a range of 0.88% to 0.9% for the example drivers shown. The longitudinal accelerating on-axis field is shown as a solid grey line (right y -axis).

electron energies W_{kin} and corresponding Lorentz factors $\gamma = 1 + W_{\text{kin}}/(m_e c^2)$, the individual electron velocities v are all very similar and close to c , due to $v = c\sqrt{1 - (1/\gamma^2)}$. Therefore, the driver beam current profile remains similar, and in turn the produced wakefield. That said, naturally, higher energy drivers are capable to drive the wake over longer distances, and hence are able to sustain the linear energy gain of the witness to higher energies.

The linear energy gain, enabled by the phase-stable nature of the PWFA, remains largely unaffected by variations in the transverse evolution of the driver beam – such as betatron oscillations and lensing effects – which are directly linked to the driver’s slice energy.

The energy gain is not entirely identical, as shown in Fig. 4.4 (b), which displays the spectral charge densities of the electron witness bunches – representing artificial energy spectra that would correspond to experimental observables. These plots reveal that the witness energy is slightly higher for higher-energy drivers compared to the baseline. They also exhibit differences in transverse divergence, reflecting an asymmetry in transverse emittance between the two projection planes. This asymmetry originates from the linear polarization of the injector laser pulse, which increases the thermal emittance [195] of the witness electrons along the polarization direction – here, the y -direction. As a result, the normalized projected emittance is $\epsilon_{n,y} = 29.0$ nm rad for the baseline witness, compared to $\epsilon_{n,x} = 12.3$ nm rad in the orthogonal direction.

These spectra are taken at the cut-off location marked by the black dashed line in Fig. 4.4 (a), representing a potential experimental extraction point designed to obtain stabilized witness energy even under extreme experimental conditions of wildly varying driver beam energies from shot-to-shot at mean witness energies exceeding the baseline driver value ($W_W > W_D$). Such a tunable extraction point could be realized using tunable-length gas cell targets [212], which today are even commercially available.

In this scenario, driver beam energies in the range of $\sim 450 - 550$ MeV yield witness beam energies with a precision better than $\pm 0.1\%$, while driver energies above ~ 360 MeV can maintain witness energy stability within $\pm 1\%$. These parameter scans demonstrate that the output energy stability of the witness beam can exceed the stability of the driver beam by more than an order of magnitude.

Plots of the longitudinal phase space for both driver and witness populations highlight effects that depend on the assumed relative energy spread of the driver beam, which here remains at the baseline 10%. The lowest-energy electrons within the driver beam that are located in its strongest decelerating region, are slowed down most rapidly. This leads to a progressive lengthening of the driver beam, which in turn reduces its peak current and weakens the wakefield strength during driver depletion.

Eventually, as shown in Fig. 4.4 (c), driver beam electrons are fully depleted, and slip into the accelerating phase of the wake. This effect, already observed in simulations [75, 213] and seen experimentally [214, 215], is naturally most pronounced for low-energy driver beams with large energy spreads. Accordingly, it is most visible in Fig. 4.4 (c) for the case of a 0.4 GeV driver. As shown by the colorbar (normalized across Fig. 4.4 and Fig. 4.5), the (projected) relative energy spread (r.m.s.) of the witness beams is $< 1.0\%$. In detail, the 0.4, 0.5, and 1.0 GeV driver beams produce witness beams with mean energies of 0.506, 0.508, and 0.510 GeV, with projected energy spreads of 0.89%, 0.90%, and 0.88%, respectively. Compared to the driver electron beam energy spread of 10%, this constitutes a reduction of more than an order of magnitude and signifies a quality enhancement of the extracted witness beam in terms of its transverse emittance and concurrently its energy spread, enabled by the hybrid LWFA→PWFA staging.

Simultaneously, the hybrid LWFA→PWFA platform therefore functions as an energy and beam quality transformer [192, 193]. The energy and energy

spread effects come with a trade-off in overall energy efficiency between the input LWFA beam and the output beam generated by the plasma photocathode in the PWFA stage. However, as the reproducibility of the LWFA-generated driver beam improves, the efficiency can be further optimized. This optimization is fundamentally governed by the depletion length of the driver beam with the lowest mean energy. As a result, a capping point closer to the maximum baseline witness energy – approximately ~ 800 MeV in this case – becomes achievable.

4.4 Witness energy gain is independent of driver beam energy spread

However, it is not only the mean energy of the driver beam that can vary from shot-to-shot. It is well-known that the energy spread of LWFA-generated electron beams can also fluctuate significantly.

To determine the effect this may have on the plasma photocathode-generated witness bunch in the PWFA stage, a scan was undertaken by varying the driver's relative uncorrelated energy spread between 0 and 50%, while keeping the mean driver electron energy constant at its baseline level of $W_D = 500$ MeV. Fig. 4.5 (a) shows the corresponding witness energy gain curves with the longitudinal phase space plots of driver and witness for selected driver energy spreads displayed in Fig. 4.5 (b).

As evident from Fig. 4.5 (a), the energy gain remains nearly unaffected by even large variations in the driver energy spread during the early stages of acceleration.

With increasing acceleration distance and driver beam deceleration, in plasma waves driven by driver beams with larger energy spreads such as 50%, the witness energy gains begin to reduce, for similar reasons as discussed in the context of varying driver beam mean energies. Notably, lower driver energy spreads allow the witness beam to reach its maximum energy earlier.

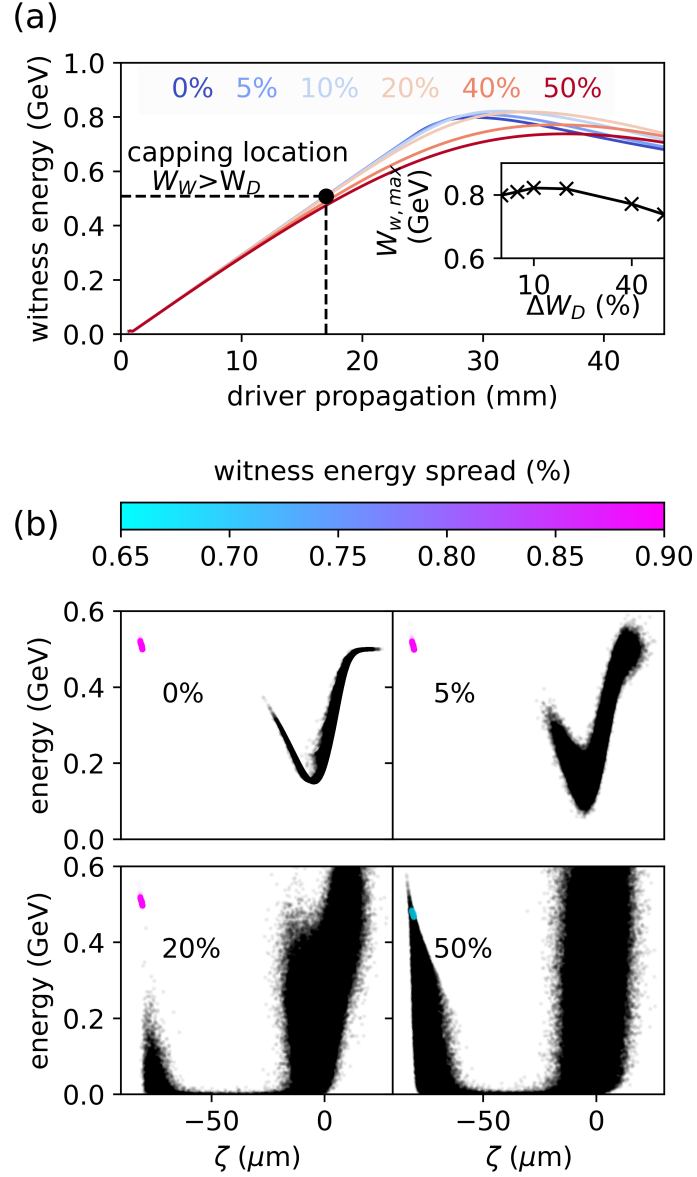


Figure 4.5: Impact of driver beam energy spread on witness bunch energy gain. Panel (a) illustrates the witness energy gain as a function of driver energy spread for a fixed driver energy of $W_D = 500$ MeV. The dashed line indicates the threshold at which the witness energy exceeds the driver energy. The inset shows the maximum mean energy gain trend for the witness bunch. Panel (b) presents longitudinal phase space plots for the driver and witness bunches at different driver energy spreads (0%, 5%, 20%, 50%), while the witness bunch maintains a relative projected energy spread of less than 1%. The plots demonstrate the influence of driver energy spread on the longitudinal distribution of both bunches.

While the largest energy spread drivers yield the lowest peak witness energies, drivers with moderate energy spread such as 10% can actually achieve slightly higher maximum energy gains than the hypothetical zero-spread case (see Fig. 4.5 (a) inset). The acceleration distances at which these maxima occur vary non-trivially, reflecting the complex interplay of wakefield dynamics and pump depletion effects.

Again, introducing a capping point – e.g., where the witness energy exceeds the driver energy – helps ensure stable witness energy output, even with strongly varying driver energy spreads from the LWFA stage. In this case, the lowest energy spread cases reach the target energy before those with higher energy spread. At the chosen extraction point, the mean witness beam energy is 507.0 MeV, and if the driver energy spread remains below 26.5% r.m.s., the extracted witness energy stays within a $\pm 1\%$ tolerance, demonstrating the hybrid LWFA→PWFA scheme’s broad tolerance to energy spread jitter. Additionally, all resulting witness bunches exhibit projected energy spreads below 1%, reducing the witness energy spread by more than an order of magnitude compared to the baseline driver beam.

In this context, it is important to note that the selected robust working point (Fig. 4.3) ensures that witness electron trapping occurs in a region of the wakefield with comparatively lower accelerating gradients, rather than at the wake’s vertex. While the vertex offers higher initial wakefield amplitudes – and thus the potential for greater witness energy gains – trapping in the chosen region provides enhanced beam stability. This placement reduces the witness electrons’ vulnerability to wakefield evolution and lowers the risk of slipping out of the focusing and/or accelerating phases of the wakefield.

Another consequence of positioning the witness bunch in a more advanced phase of the wakefield is the potential for spatial overlap with decelerated electrons from the driver beam. Here, they may contribute to beam-loading [216],

which may be exploited for an interesting effect in particular in context of ultralow emittance electron beams. In Ref [209], it has been found that beam-loading by an ‘escort’ beam released to overlap with an ultralow emittance witness beam that has already been accelerated to sufficiently relativistic energies, can reduce or even compensate the witness energy chirp, without spoiling its emittance. In the present configuration, such beam-loading can be provided by the decelerated drive beam fraction that slips back into the accelerating phase of the wakefield. Fig. 4.5 (b) shows the signature of this effect in the snapshot for the scenario of the 50% driver energy spread. Here, the witness’ projected energy spread is significantly lower ($\sim 0.68\%$ vs. $\sim 0.9\%$) than in the other three snapshots shown, as the result of beam-loading (see Section 4.5). For completeness, the 0%, 5%, 20% and 50% driver energy-spread cases yield mean energies of 0.509, 0.508, 0.506 and 0.474 GeV, and corresponding projected relative energy spreads of 0.90%, 0.90%, 0.89% and 0.68%. This creates an opportunity to purposefully reduce the witness beam’s energy spread while increasing its energy, provided there is sufficient control over the driver beam. While this scenario requires further exploration, it could be a natural and perhaps rather unique fit to hybrid LWFA→PWFA-plasma photocathodes, since they can provide sufficiently large energy spread driver beams, and ultralow emittance witness beams sitting in sufficiently advanced phases of the wakefield for spatial overlap with decelerated driver beam electrons. Improving the energy spread while keeping the low emittance could be crucial for applications such as free-electron-lasers, which impose strict thresholds on both parameters [118, 217].

4.5 Beam loading and witness chirp reduction by driver beam depletion

As discussed in Section 4.4, certain latent effects which have an observable impact on the witness beam's energy spread become noticeable at longer acceleration distances. Fig. 4.6 shows the underlying cause of the reduced energy spread behaviour observed in the energy phase space diagrams in Fig. 4.5 (b).

The top three panels show the longitudinal driver and witness current profiles at the capping point, and the longitudinal wakefield profiles generated by a driver beam of 5%, 10% and 50% energy spread at injection (black solid line) and at the capping point (dashed line). The initial driver energy is 500 MeV and its charge 500 pC.

While there is no difference between the wakefields at the beginning of the acceleration and the end for the low energy spread cases, the 50% energy spread case shows that there is significant difference in wakefield amplitude at the witness position. As the driver beam depletes due to the self-created decelerating phase of the field, it causes the lowest energy electrons to eventually be recaptured at a phase of the wakefield that overlaps with the witness beam. This additional, significant current at witness position provides local beam loading, and flattening of the longitudinal wakefield at this position. The bottom three panels are zoom-ins, highlighting the change of the wakefield due to this decelerated and re-accelerated charge. This flattened wakefield is responsible for the reduction in witness energy spread.

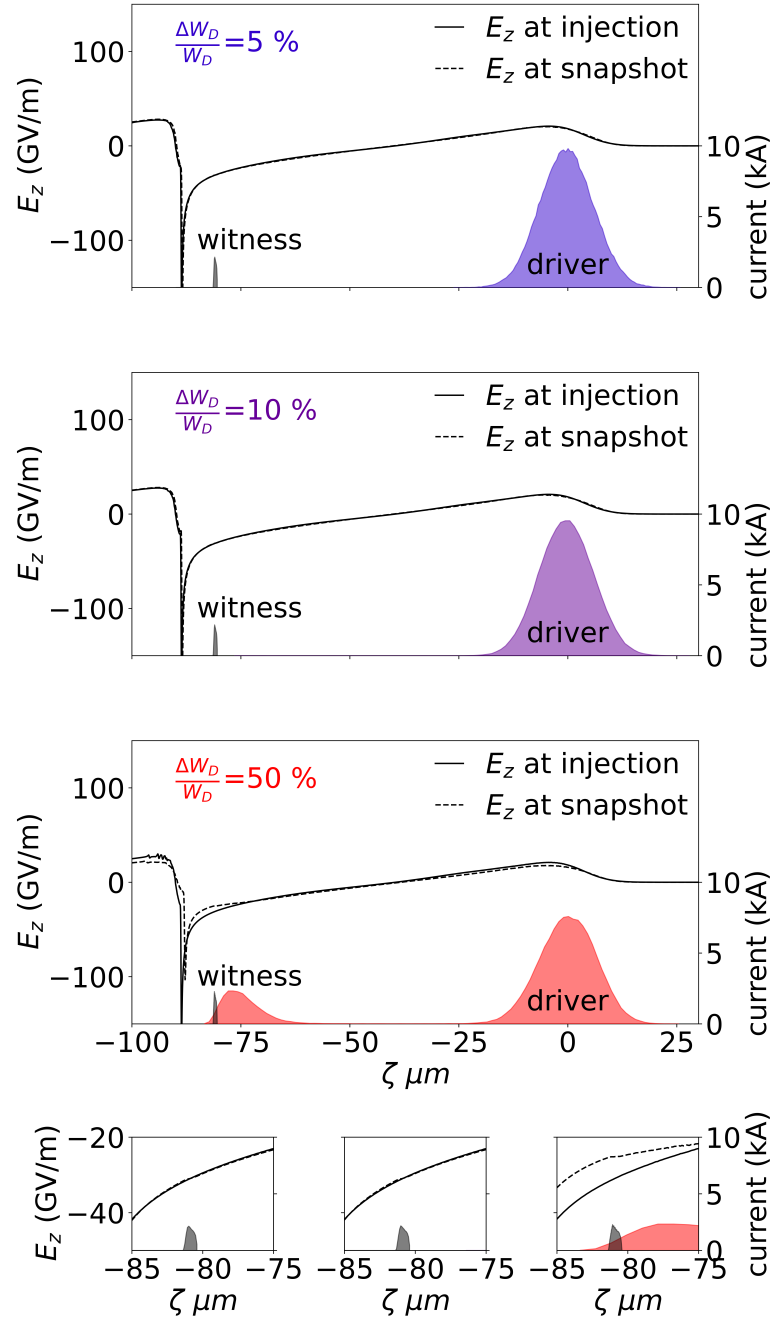


Figure 4.6: Wakefield evolution showing beam-loading effects from driver energy spread. Top panels: On-axis accelerating field at injection (solid black) and at 17.0 mm (dashed black), shown with color-coded driver current profiles and shaded black witness current. Increased driver energy spread accelerates depletion, causing field flattening around the witness beam position. Bottom panels: Detailed view of the flattened accelerating field region across the witness bunch, demonstrating how driver depletion contributes to reduced energy spread in the witness beam.

4.6 Witness production has high spatio-temporal release tolerance

Next, the influence of the spatio-temporal release position of electrons by the plasma photocathode within the plasma wave is examined. As before, the baseline driver beam is assumed. A test-particle approach is employed to simulate all possible electron release positions simultaneously, independent of specific laser parameters. The results show that electrons released within the solid grey ellipse (see Fig. 4.7) are captured with 100% capture efficiency, accelerated, and retained throughout the entire acceleration length.

As expected, this aligns with the analytical solution (dashed grey) for trapping from negligible initial velocity, corresponding to the region satisfying the simplified trapping condition discussed earlier in Section 4.1, here applied in 2D.

Application-specific witness bunch energy stability targets then inform the acceptable spatiotemporal plasma-photocathode release region in the wake as a subset of the trapping ellipse. Iso-contour lines show where released, and subsequently trapped and accelerated electrons would obtain 0.1% (yellow), 1.0% (red), 5.0% (blue), 10.0% (black) energy deviation with respect to electron release at electrostatic potential minimum. The symmetry of the injection acceptance energy iso-contours reflects the underlying symmetry of the electrostatic potential around the centre of the wakefield. In hybrid LWFA→PWFA configurations, the temporal synchronization jitter of the injection laser may be less critical than its pointing jitter. To achieve the 1% witness energy stability target considered in this study, the transverse shot-to-shot pointing fluctuation must be limited to below $\pm 3 \mu\text{m}$ —a level of stability within reach of emerging high-power laser system capabilities [189, 218, 219].

The positive gradient of E_z , in concert with linear focusing fields, E_r , facilitates automatic longitudinal and transverse compression of electrons at the end

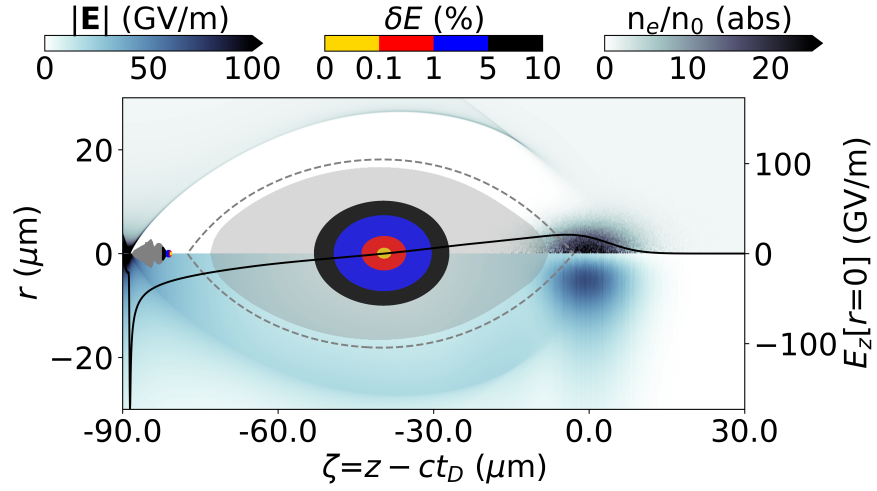


Figure 4.7: Intrinsic robustness of witness bunch production against photocathode laser spatio-temporal jitter, shown as a function of longitudinal and radial release position within the blowout. Plasma and driver electron densities are shown in a white-light blue-green gradient, and the localized electric field magnitude is depicted in a white-dark blue gradient. Electrons released within the color-coded energy bands around the 2D trapping ellipse (grey area) are trapped in small volumes at the back of the blowout. The dashed grey line represents the 1D theoretical trapping ellipse, with energy deviations (δE) marked by contour boundaries and the on-axis longitudinal electric field shown as a solid black line.

of the plasma wave bucket. The real space volumes of trapped test electrons corresponding to the release volumes are indicated with the same colors as the target iso-contour lines in Fig. 4.7. The auto-compression is a particularly beneficial feature of the aforementioned injection and trapping process, automatically generating witness bunches with ultrashort bunch durations and multi-kA peak currents, exceeding 2.2 kA in the showcase configuration.

When injection occurs off-axis, betatron oscillations are induced by the focusing fields; however, these oscillations are rapidly damped upon acceleration given that the electron beam’s relativistic mass increases with γ , and so the same focusing force produces less transverse acceleration. Effectively, the witness bunch follows the orbit of the driver beam which leads to smaller emittance in comparison with external injection schemes involving already relativistic witness beams, where even minor transverse offsets or pointing errors can severely compromise the acceleration process [208]. For instance, recent studies have indicated [220] that already micrometre-scale emittance growth becomes difficult to control with even sub-micrometre misalignment. In contrast, simulation studies suggest that plasma photocathodes can limit emittance growth to the nanometre level – more than two orders of magnitude lower – even in the presence of substantially larger transverse offsets [169]. It should be highlighted that these simulations do include the effects of exiting the plasma, utilising a so-called ‘escort’ beam to reverse the witness beam’s energy chirp and create the low projected energy spread and flat longitudinal phase space at the end of the plasma stage required for nm rad-level emittance preservation.

4.7 Witness beam energy gain self-adapts to driver charge density fluctuations

In Section 4.2, the driver charge density was varied to demonstrate the witness beam injection rates in the centre of the wake are independent of the driver charge density, and the witness charge yield is hence fully decoupled from shot-to-shot-variations of the driver beam. However, as is well-known from theory [200, 221, 222], experiments and even directly observable via shadowgraphy [79], variations in the driver density relative to the plasma density n_D/n_p due to driver charge Q_D , length σ_z , transverse size σ_r or shape variations from shot-to-shot, can also strongly vary the peak wakefield amplitude, the plasma wavelength and hence the plasma wave cavity size. Therefore, while the accelerating wakefield is only marginally dependent on driver beam energy and its energy spread as discussed in Sections 4.3 and 4.4, in contrast it is strongly dependent on driver charge density variations [80].

It is also well-known that injection rates in LWFA are difficult to control [98, 223], and the slice and total current values of electron beams produced by LWFA therefore often exhibit even larger shot-to-shot variations than in energy or energy spread. While total charge is easy to measure, variations in electron beam length and shape are notoriously difficult to quantify. Consequently, shot-to-shot variations in LWFA-produced driver beam charge density represent a very significant, yet often elusive and challenging-to-measure, source of variation for wakefields in hybrid LWFA→PWFA, ultimately influencing the energy gain experienced by witness beams injected in the PWFA stage.

In the hybrid LWFA→PWFA scenario, it is realistic to assume charge release occurs at a fixed distance behind the driver beam, due to the intrinsic synchronization with the injector laser pulse, so how then does the witness energy gain change in this scenario as consequence of driver beam density jitter. To inves-

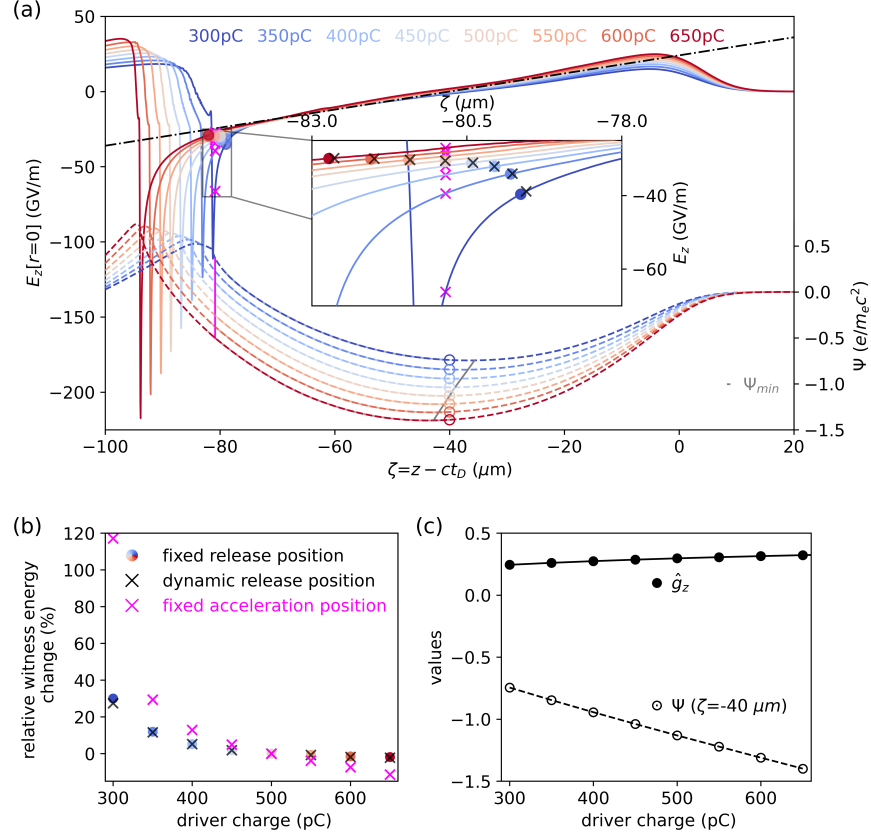


Figure 4.8: Intrinsic mitigation of witness energy fluctuation arising from driver density-induced wakefield jitter. Panel (a) presents wakefield line-outs for various driver charges, highlighting changes in the electrostatic potential and shifts in the location and magnitude of the minima for both E_z and $\Psi_{\min}$. The dash-dot black line denotes the theoretical normalized gradient $\hat{g}_z$ of the reference driver. Overlaid are color-coded hollow circles indicating fixed injector release positions, with corresponding trapping locations marked by solid circles of the same color. Black crosses represent injection at local $\Psi_{\min}$ positions, while magenta crosses correspond to injection at fixed trapping locations. Panel (b) compares the relative changes in witness energy gain: releasing electrons at a fixed location near the blowout centre closely matches injection at the local $\Psi_{\min}$, whereas the variation of mean energies is significantly larger when the acceleration location would be fixed. Panel (c) shows Ψ assuming a fixed injector release position at $\zeta = 40$ μm (hollow black circles), and the normalized slope $\hat{g}_z$ (solid black circles) approaching an asymptotic value of $1/2$, as a function of driver charge.

tigate this, the drive beam density was varied as in section 4.2, and the plasma photocathode release position was set to the electrostatic minimum for the case of the driver charge of $Q_D = 500$ pC. The variation of charge from $Q_D = 300$ pC to 650 pC corresponds to a relative change of -40% to +30%.

Fig. 4.8 (a) visualizes the significant variation in electrostatic potentials and longitudinal wakefield configurations as a response to changes in driver charge (density). The color-coding emphasizes that wakefields are stronger (red) and weaker (blue) for driver charge higher/lower than the baseline driver charge $Q_D = 500$ pC. The color-coded hollow circles represent the plasma photocathode release position ζ_{rls} , and highlight that on the one hand the electron release positions are located at electrostatic potentials that are significantly different for the range of Q_D considered. On the other hand, for all cases other than the baseline driver charge, they are also located not at the electrostatic potential minimum of the waves, but to the right (red, stronger waves) and left (blue, weaker waves) of their individual minimum, represented by the dashed grey dotted line. However, the positions in the wakefield where the electrons are trapped (ζ_{trap} , semi-transparent solid circles) are closer to the release position for weaker waves (blue) than for stronger waves (red). This may be counter-intuitive, because in a constant wakefield, release positions outside the electrostatic potential minimum result in trapping positions ζ_{trap} farther behind in the plasma wave. In phase-constant PWFA, the wakefield at the trapping position ζ_{trap} then determines the energy gain.

These findings highlight an intrinsic stabilization mechanism inherent to the plasma photocathode injection scheme, rendering it particularly advantageous in the context of charge variability characteristic of hybrid LWFA→PWFA systems.

The inset in Fig. 4.8 (a) zooms in on the trapping region, highlighting that compared to a hypothetical fixed acceleration position (magenta crosses), for comparison taken here to be the same as the trapping position resulting from

plasma photocathode release for the baseline driver case, the range of electric wakefields at the individual trapping positions ζ_{trap} is significantly reduced across the range of driver beam strengths. This represents a self-adaptive mechanism that decreases the energy gain fluctuations of the witness beam despite significant fluctuations in drive beam and wakefield amplitudes.

In other words, while weaker waves have smaller electrostatic potential and peak wakefield amplitudes than stronger waves, at fixed wake phase positions, the accelerating field – and hence the acceleration rate – is significantly larger in weaker waves than in stronger ones. This is a challenge for the output energy stability of externally pre-accelerated and then injected electron beams, in addition to timing jitter. While in the plasma photocathode mechanism, the resulting accelerating wakefields $E_{z,\text{trap}}$ at individual trapping positions ζ_{trap} are also larger for weaker waves than for stronger ones, the difference is significantly mitigated: trapping positions ζ_{trap} automatically shift to later positions within the wake for stronger waves – where $E_{z,\text{trap}}$ is larger than at fixed trapping position. This dynamic shift partially compensates for variations in driver beam strengths from shot to shot, thus providing intrinsic mitigation of the actual witness acceleration gradient fluctuation, and hence the witness energy gain fluctuation. In further detail, the inset in Fig. 4.8 (a) shows that as the driver charge is linearly increased, the variation of ζ_{trap} increases, while the change in $E_{z,\text{trap}}$ decreases non-linearly. This leads to the bow-shaped distribution of wakefield vs. trapping position displayed in the inset.

Fig. 4.8 (b) quantifies the witness energy fluctuation mitigation by comparing the resulting witness energy change obtained for beams released by the plasma photocathode with those that would be obtained at a fixed acceleration position in the wakefields. The y -axis shows the relative change compared to the baseline driver case of $Q_D = 500$ pC. The magenta plot represents the hypothetical fixed wakefield position case.

The colored circles show the electric fields at trapping positions resulting from plasma photocathode emission at a fixed release position, and the black crosses show the slightly different electric fields at trapping positions when release occurs at the individual potential minima of each driver-generated plasma wave. This comparison highlights the intrinsic energy stability advantage of electron beams that can be expected from plasma photocathode injectors.

The - to some extent - self-stabilizing behaviour of the witness energy – and its robustness to increasing driver current, which leads to saturation of the energy gain in the high-current limit – can be understood from first principles. Consider an accelerating electric field with constant slope, $g_z \equiv dE_z/d\zeta > 0$, defined along the co-moving coordinate $\zeta = z - ct_D$ and assume that the electron is released at the zero crossing of the parent wakefield, ζ_0 , corresponding to the centre of the blowout region where $E_z(\zeta_0) = 0$. The field is then given by

$$E_z(\zeta) = g_z (\zeta - \zeta_0). \quad (4.1)$$

Applying the trapping condition, $\Psi(\zeta_0) - \Psi(\zeta_f) = -mc^2/e$, and using $E_z = \partial_\zeta \Psi$, integration yields the trapping offset

$$\zeta_f - \zeta_0 = -\sqrt{\frac{2m_e c^2}{e g_z}}. \quad (4.2)$$

The accelerating field at the trapping location then evaluates to

$$E_{z,} = -E_0 \sqrt{2\hat{g}_z}, \quad (4.3)$$

where $E_0 \equiv m_e c^2 k_p / e$ is the cold wavebreaking field, and $\hat{g}_z \equiv g_z / (E_0 k_p)$ is the normalized slope of the wakefield. This result shows that the final accelerating gradient depends solely on $\hat{g}_z$.

As established by theory and simulations [200, 221], the parameter $\hat{g}_z$ asymptotically approaches 1/2 for high-current drivers ($\gg 10$ kA). In the regime of moderate currents (around ~ 10 kA), $\hat{g}_z$ exhibits only a weak dependence on

the driver's current, typically settling near 0.3, in agreement with earlier findings [224]. Due to the square-root dependence in Eq. 4.3, the influence of current on $E_{z,\text{trap}}$ is further diminished, resulting in a slow convergence toward $E_{z,\text{trap}} \rightarrow -E_0$ as the beam current increases.

This behaviour is further illustrated in Fig. 4.8 (c): the difference of Ψ at plasma photocathode release position of $\zeta = 40 \mu\text{m}$ (hollow circles) shows a strong dependence on driver charge, whereas $\hat{g}_z$ (solid circles) – and by extension also $E_{z,\text{trap}}$ – converges towards increasing driver charges. For lower driver currents, trapping occurs increasingly deeper within the non-linear blowout region, where deviations from linear theory become significant. Nevertheless, the injected beam remains more stable compared to trapping at a fixed location, further illustrating the resilience of this injection mechanism.

4.8 Summary

The plasma photocathode is a path towards electron beams with ultralow emittance ϵ_n of the order of tens of nm rad in both planes, kA-level current I and in turn ultrahigh electron beam brightness $B_{5,n} = 2I/(\epsilon_{x,n}\epsilon_{y,n})$. While not discussed above, also in the present study the witness beam brightness exceeds the driver beam brightness by orders of magnitude, and hence acts as a brightness transformer for PWFA. The projected 5D-brightness consistently reaches values of $B_{5,n} > 10^{19} \text{A m}^{-2} \text{rad}^{-2}$, and projected 6D-brightness of the order of $B_{6,n} > 10^{18} \text{A m}^{-2} \text{rad}^{-2}/0.1\% \text{BW}$.

Both the ultralow emittance and ultrahigh brightness potential of plasma photocathodes make them ideal candidates for photon sources such as soft X-ray coherent synchrotron radiation sources [225], hard X-ray FEL [118], betatron radiation and ion channel lasers, and γ -ray sources [226]. Such beams could also be exploited for R&D toward beam quality of high energy physics (HEP) collid-

ers and for next-generation experiments probing quantum electrodynamics. The paramount importance of beam quality and in particular emittance and brightness for both thrusts is well-known [227], and hence, for example, plasma-based X-FELs are seen as a major milestone and stepping stone toward high energy physics applications [228]. As discussed in Section 4.6, techniques are being established that will allow beam brightness to be maintained during extraction from the plasma stage [209].

LWFA driven FELs have previously been demonstrated [229], however as FEL performance depends strongly on beam properties such as energy spread, emittance and overall stability, achieving stable FEL gain is challenging, especially for X-rays, due to shot-to-shot fluctuations in the laser being imparted on the witness electron beam. PWFA FELs have also been demonstrated [230] and are generally closer to conventional FEL requirements for hard X-ray production due to this method's ability to produce ultra-low emittance and high brightness beams. The hybrid LPWFA concept discussed in this chapter combines the compactness of an LWFA setup with PWFA beam quality and stability, offering a more promising route to practical plasma-based FELs.

In such schemes, driver and witness beams differ in energy, emittance and divergence, which allows the separation of the driver in order to isolate the witness beam in a chicane line via a permanent quadrupole magnet (PMQ) triplet placed at a defined drift distance from the end of the plasma to capture the witness bunch. Constructing the chicane to have a small bend angle $\theta \approx 2$ mrad changes the orbit of the high-energy witness beam negligibly, thus avoiding witness beam quality loss by coherent synchrotron radiation. In contrast, driver electrons are kicked out several millimetres and can thus be dumped, or exploited for diagnostics or other applications.

Current hybrid LPWFA plasma photocathode setups have demonstrated witness beam charge of $Q \approx 8.5$ pC [231] which falls just below the range of conven-

tional RF linac FELs $Q \approx 20 - 250$ pC [232]. However successful FEL operation is more directly dependent on peak current, with hard X-ray FELs typically requiring $\approx 2-4$ kA, energy spread, with FEL lasing in the X-ray regime requiring $\Delta W/W \lesssim 0.1-1\%$, and emittance, with $\approx 0.2-0.5 \mu\text{m}$ rad needed for hard X-ray FELs, enabling plasma photocathode beams to satisfy the necessary criteria.

Although emittance, brightness and energy spread are fundamentally decisive parameters for applications such as light sources, the stability and reproducibility of electron beams are also important. Plasma photocathodes can be realized with driver beams from conventional RF accelerators [116], but as of today very few linacs offer PWFA capabilities, in particular high-field PWFA capabilities with plasma wakefields > 10 GV/m. Only one linac engaged in PWFA research – SLAC FACET [116, 233] – so far offered a capability crucial for hosting plasma photocathodes, namely producing a driver beam current $\gtrsim 6$ kA.

Fortunately, LWFAs can produce beams with very large currents, allowing hybrid LWFA→PWFA systems to even operate in the high-field PWFA regime, and to mature from concept [75] to experimental reality in recent years [76–80]. These systems were quickly envisioned as platforms capable of hosting plasma photocathodes, with their intrinsic synchronization capability recognized as an important additional advantage [93]. Very recently, the first plasma photocathode in a hybrid LWFA→PWFA system was realized, at HZDR’s DRACO facility, utilizing a 90 degree geometry [231].

With these advances, and the global number of high-power laser installations now reaching into the triple digits, the availability of LWFA-capable facilities that can be leveraged for PWFA and plasma photocathode research has grown substantially. This makes it both timely and necessary to investigate the specific conditions and parameters governing this hybrid plasma photocathode scenario in greater depth.

In this chapter, the principal sources of input jitter anticipated when im-

plementing an LWFA→PWFA equipped with a collinear plasma photocathode setup were both identified and explored. The dominant shot-to-shot variations of plasma waves driven by electron beams from today’s LWFA systems originate from fluctuations in the electron beam’s energy, energy spread, and charge or current. Then, plasma photocathode injection jitters in these plasma waves may result from spatio-temporal jitter between the plasma photocathode laser and the electron beam driven plasma wave, and from shot-to-shot changes in the laser intensity.

The findings presented show that plasma photocathode electron injection rates are purely dependent on the injection laser, and the witness electron beam charge therefore is decoupled from the driver beam strength shot-to-shot. This is in stark contrast to injection methods such as downramp injection or drive beam- or wakefield-induced ionization based methods, where injection rates are strongly dependent on driver and wakefield strength.

Regarding the acceleration gradient, fluctuations in the driver beam energy and energy spread exhibit negligible impact until driver depletion occurs. Lower-energy driver beams deplete earlier; thus, limiting the length of the PWFA stage can stabilize the output witness energy and operational reproducibility, as alternative to higher witness beam energies, but reduced reproducibility. Interestingly, driver beam depletion especially in case of large driver energy spreads can induce beam-loading in the witness beam trapping region, which can lead to a significant reduction of witness energy spread.

Targeting the centre of the blowout for plasma photocathode release benefits from the symmetry and flatness of the electrostatic wake potential in this region. This symmetry provides intrinsic robustness against shot-to-shot spatio-temporal relative jitter, as small variations in the electrostatic potential at the release position result in only minimal shifts in the trapping position. Consequently, the associated electric fields and accelerating gradients at the trapping position

remain similar shot-by-shot.

It shall be emphasized that an injection location in the centre of the wave is a feature uniquely available for plasma photocathode injection. In contrast, other internal injection methods typically either inject in earlier (e.g. driver pulse-induced ionization injection) or in later (wakefield-induced or density downramp injection) wake phases, resulting in trapping positions further towards the rear of the wake. In contrast, plasma photocathodes naturally enable earlier trapping positions within the plasma wave, offering significant advantages on multiple fronts. First, the wakefields at the release position are minimal, ensuring that their contribution to the witness charge yield due to superposition with the plasma photocathode laser fields is negligible. Second, witness beams are trapped and accelerated in a ‘safe zone’ sufficiently well separated from the vertex of the wake. On the one hand, this allows robust, full capture of the released charge, as well as avoiding witness beam vulnerability to longitudinal and transverse wake evolution and potentially even charge loss. On the other hand, the wakefield amplitude differences at neighbouring wake positions rather deep into the blowout are rather small, minimizing the spread of accelerating gradients in case of jitters. Third, it is revealed that in combination with the intrinsic synchronization of hybrid systems and the minimized temporal jitter between plasma photocathode and driver beam, an automatic significant stabilization mechanism against driver charge variation can be harnessed. At lower driver beam charges the electrostatic wake potential is smaller than at higher driver beam charges; however, the concurrent shift in trapping positions significantly alleviates the spread of accelerating gradients of the witness beams – an effect that is favourable for witness beam energy stability. Plasma density variations are not considered in this study. A brief sensitivity analysis of plasma density and entrance ramp-length variations indicates that, within realistic experimental fluctuations, these effects induce only marginal changes in the wakefield structure and witness beam parameters.

In summary, the work presented in this chapter shows that counter-intuitively, the high beam quality feature of plasma-photocathodes does not come at the price of reduced stability, but is concomitant with improved stability. It is revealed that the plasma photocathode can act not only as a brightness and electron energy transformer, but also as a stability transformer. These benefits are made as a trade-off with beam charge. The scenario of intrinsic synchronization but significantly varying driver beam from shot-to-shot, the reality of today's LWFA systems, is complementary to the scenario of limited synchronization but state-of-the-art driver beam stability in linac-driven PWFA systems.

The plasma photocathode thus emerges as an ideal injection mechanism for hybrid LWFA→PWFA systems, offering a unique combination of three key performance enhancements – brightness, energy, and stability – delivered simultaneously. The concurrent optimization of these beam qualities is essential for demanding applications such as driving free-electron lasers, thus making this avenue an attractive route towards compact, yet high-performance accelerator systems.

Chapter 5

Plasma afterglow metrology

Regularly monitoring relevant laser and gas parameters is essential for maintaining sustained high quality performance from a laser-plasma accelerator, as fluctuations to one or more of these parameters can significantly alter the interaction physics and even prevent electron acceleration all together. As detailed in Chapter 4, ionization-based injection methods are being adopted that rely on auxiliary laser pulses for additional plasma generation purposes [116, 117, 234]. Variability in this secondary pulse’s intensity can therefore lead to randomness in the injected beam current, increased normalised emittance and ultimately reduced beam quality. Many plasma accelerators, like the kHz LPA described in Chapter 3, are also reliant on a tailored gas or electron density profile to gain control over the injection process and keep it localized so that it is separate from the acceleration stage [94, 235]. The ramps of the gas profile are of particular importance due to plasma focusing [236] e.g. in the context of staging multiple plasma accelerator modules [237] and for emittance preservation [238]. Therefore, variations in the gas flow can lead to a similar outcome in terms of non-optimised electron beam parameters that are unsuitable for the selected application.

At present, measuring ultrashort laser pulse parameters requires a collection of diagnostics being set up independently and at different positions along the

experimental beamline [239–241], many of which are highly complex such as an optical probe. Additionally, material damage thresholds prevent intercepting devices for more concentrated and intense beams of light, meaning these apparatus often have to be placed far away from the interaction point (IP) and operate at low laser powers. However, these factors inherently decrease the accuracy of the measurement and can conceal key details about the laser’s evolution inside the gas or plasma. This is particularly relevant when we consider typical laser-plasma acceleration regimes, where the onset of complex non-linear mechanisms such as self-focusing and/or ionization-induced defocusing are critical to the dynamics/evolution of the acceleration process. Furthermore, it is particularly important to note that many of these diagnostics are destructive, meaning the accelerator cannot be operated in tandem with carrying out the diagnosis.

Techniques for measuring the gas properties vary depending on the form the gas is presented in. In a so called ‘flooded chamber’ or gas cell, the gas density can be uncovered by diagnosing the pressure using a suitable pressure gauge connected to a digital controller. A typical high performance gauge can give out high accuracy readings but is limited to a select range of pressure values over which they are receptive. When operating in a gas jet configuration, it is desirable to obtain the profile of the gas flow [242] which is typically identified using some form of interferometry [243], a technique which generally relies upon analysing the phase shift induced by a change in the index of refraction due to the density of the gas. Unfortunately, this becomes increasingly unreliable as you move towards lower gas densities as the phase shift imparted on the incident probe beam can become undetectable.

Many of the aforementioned obstacles can however be addressed by employing a flexible optical metrology approach - plasma ‘afterglow’ (see Section 2.3.3). This technique was discussed in the context of plasma accelerators in [244] where the authors established it as a way of measuring the driver electron beam energy

loss due to ionization on a shot to shot basis during a PWFA experiment at the Stanford Linear Accelerator Center (SLAC). These capabilities were then expanded upon at the Facility of Experimental Tests (FACET) [47], where the observable afterglow was exploited to establish the concurrence, both spatially and temporally, between a relativistic electron beam and an ionizing laser pulse. More recently at the Deutsches Elektronen-Synchrotron (DESY), it was used first in the same context as FACET [122], and then to provide longitudinally resolved measurements of the energy-transfer efficiency from the plasma wake to the witness electron bunch [245]. All of these experiments chose to record and analyse the integrated afterglow count enhancement on a CCD as a function of spatio-temporal overlap or more/less energy deposition. To extend the usability of the technique beyond this, more detailed knowledge of the signal is required.

In order to establish the afterglow’s full capacity as a diagnostic method, it is crucial to map out the signal’s dependency on key experimental parameters like the laser intensity or gas density. Obtaining this knowledge will allow us to better understand the diagnostic’s behaviour in future experiments and hopefully unlock more novel use cases for it. However, this has proved to be challenging to date due to inherent fluctuations in the seed laser being imparted on the plasma filament and subsequent afterglow. This makes it difficult to alter individual parameters and confidently establish trends as these intrinsic variations in the afterglow signal could be distorting the true relationship. What is required therefore, is an extremely stable laser that permits tight control of its parameters and can provide a consistent plasma seed.

In this chapter, a highly stable Coherent Legend Elite Duo He+ laser is used to satisfy these conditions. This ‘turnkey’ system is both user friendly and boasts stellar beam pointing ($< 10 \mu\text{rad}$) and power stability ($< 0.5\%$), the latter being a by-product of its kHz repetition rate enabling a thermal steady state to be achieved. As a result, jitter contributions to the plasma afterglow were suc-

cessfully minimized. This coupled with the other inherent advantages of utilising a kHz system such as the ability to integrate over many shots and drastically improve the signal to noise ratio enabled the author to create a controlled experimental environment where trends between the afterglow and relevant experimental variables could be comprehensively studied and categorised. Furthermore, the in-situ nature of this technique allowed the author to measure the effect of ionization-induced defocusing on a high intensity laser pulse as it propagated through a gaseous medium at various densities. This is crucial for future experiments in order to properly denote the complicated interplay taking place and more accurately predict any subsequent electron generation and acceleration.

The following chapter begins by comprehensively outlining the experimental setup and methodology, starting with the laser system before moving on to the beam transport, vacuum chamber, the laser-gas interaction that creates the plasma filament and subsequent afterglow, and finally the associated afterglow diagnostic devices. Then, experimental results recorded whilst operating in flooded chamber mode are presented, beginning with a discussion on observations made of the generated afterglow signal, before delving into the dedicated parameter scans of the laser pointing angle, laser energy and gas density. Supplementary to the experimental work, simulation studies are then shown that aimed to replicate the real-life conditions of the experiment and thus recreate and better understand the afterglow's relationship to the laser-gas interaction taking place during the parameter scans. Finally, experimental results recorded whilst operating in gas jet mode are displayed that demonstrate the afterglow's capability to successfully reconstruct the profile of a gas jet nozzle at various backing pressures.

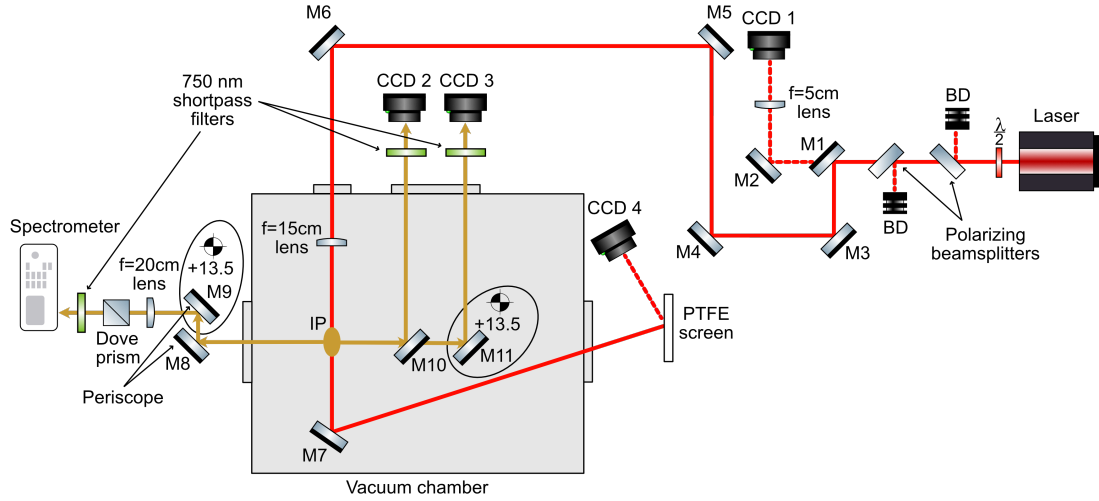


Figure 5.1: Experimental layout. Experimental setup in the SCAPA kHz lab. The laser pulse propagates through a $\lambda/2$ waveplate and two polarising beam splitters, which in tandem act as an attenuator with two beam dumps (BD) capturing the unwanted portions of the beam, before being directed towards the window port on the far-side of the vacuum chamber by a series of mirrors. A small portion of the pulse is transmitted through M1 and imaged onto CCD 1 to create an online laser energy and near field profile diagnostic. The laser passes through a 6.5 cm diameter calcium fluoride (CaF_2) window into the vacuum chamber, which is filled with Helium gas, where an achromatic doublet lens focuses it to produce a thin plasma filament at the interaction point (IP). The laser is then directed out of the chamber and onto a PTFE screen imaged by CCD 4, allowing the pulse profile and energy to be captured both pre and post interaction. The plasma afterglow radiating from the filament is observed from the side by SideView CCD 2, from the top by TopView CCD 3 and by an imaging spectrograph equipped with a dove prism that enabled the afterglow signal to be spatially spectrally resolved along both z (laser propagation axis) and y (towards the ground and ceiling). Each afterglow camera is equipped with a 750 nm shortpass filter to block any light emanating from the laser pulse. The elevation target symbol and associated number assigned to both the CCD 3 and spectrometer imaging systems signifies an increase in height of 13.5 cm.

5.1 Experimental setup

The schematic displayed in Figure 5.1 is representative of the optical beamline constructed within the SCAPA kHz lab. CCDs 1 - 4 were Teledyne FLIR Blackfly S cameras and the spatial calibration for each one was determined by inserting an object of known dimensions into the imaging system, typically a measuring tape, and observing the number of pixels that the object encompassed. Information pertaining to the resolution, filtering and orientation of these cameras is located in Table 5.1.

CCD	Vertical spatial resolution (mm/px)	Horizontal spatial resolution (mm/px)	Image dimensions (mm)	ND filters	Shortpass filters	Orientation - laser propagation
1	0.028	0.028	20.8 x 20.8	OD9	None	Into camera
2	0.029	0.032	4.8 x 4.4	None - OD0.2	750 nm OD4	Left to right
3	0.039	0.038	4.6 x 4.7	None	750 nm OD4	Left to right
4	0.114	0.034	24.0 x 123.0	OD1	None	Into PTFE screen

Table 5.1: Camera specifications. A list of the relevant settings for each CCD, including the spatial calibration, number and type of filters attached and its orientation w.r.t. the laser propagation axis.

5.1.1 Laser system characterisation

The Coherent Legend Elite Duo He+ laser employed in the experiment is an ultrafast Ti:Sa amplifier that calls upon chirped pulse amplification (CPA) to create laser pulses with $M^2 = 1.35$, capable of achieving terawatt-level peak power with $> 1000 : 1$ pre-pulse and $> 100 : 1$ post-pulse contrast ratio. More

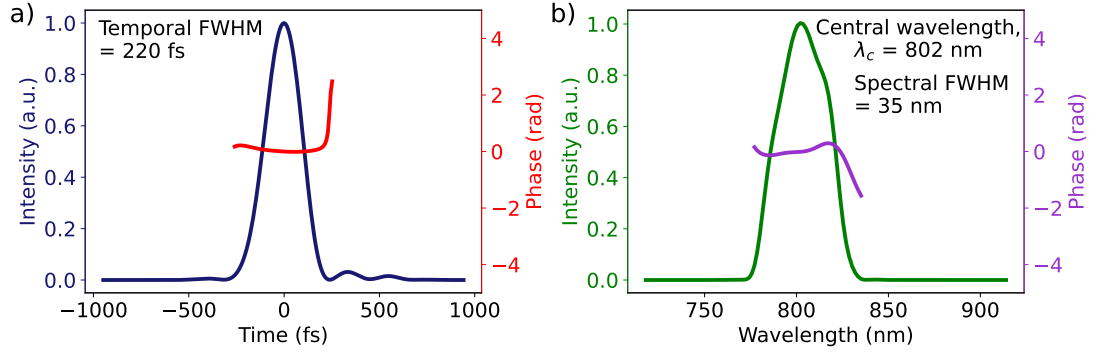


Figure 5.2: Laser pulse duration and spectrum. Measurements of both the intensity and phase of the laser pulse's a) duration and b) spectrum by the FROG device.

specifically, the system can deliver up to 13 mJ/pulse in a pulse range of 25 fs - 1 ps at a repetition rate of 1 kHz. The centre wavelength lies between 780 - 820 nm and the nominal beam diameter is 12 mm at a divergence angle of $\theta \approx 1 - 2$ mrad.

Both the spectrum and pulse duration, which was set to $t_{FWHM} \approx 220$ fs so as to be unaffected by temporal spreading as it passed through the chamber window (see Section 5.1.2), were obtained by directing the pulse exiting the laser into a frequency-resolved optical gating (FROG) device [241] and the recorded data is displayed in Figure 5.2. The FROG provides spectrally resolved readings so it was possible to retrieve both the intensity and phase of the pulse versus time and wavelength. The laser pulse's phase in time and wavelength are complementary and related via Fourier Transform, with the temporal phase describing how the instantaneous frequency evolves during the pulse and the spectral phase characterising how different wavelengths are delayed and interfere to form the pulse. The phase is measured in angles, typically radians, because it represents a position within a cycle of a wave.

The laser's near field profile is shown in Figure 5.3, which was obtained by taking the small fraction of laser light that is transmitted by mirror M1, further attenuating it using several ND filters and then imaging it onto CCD 1 using a f

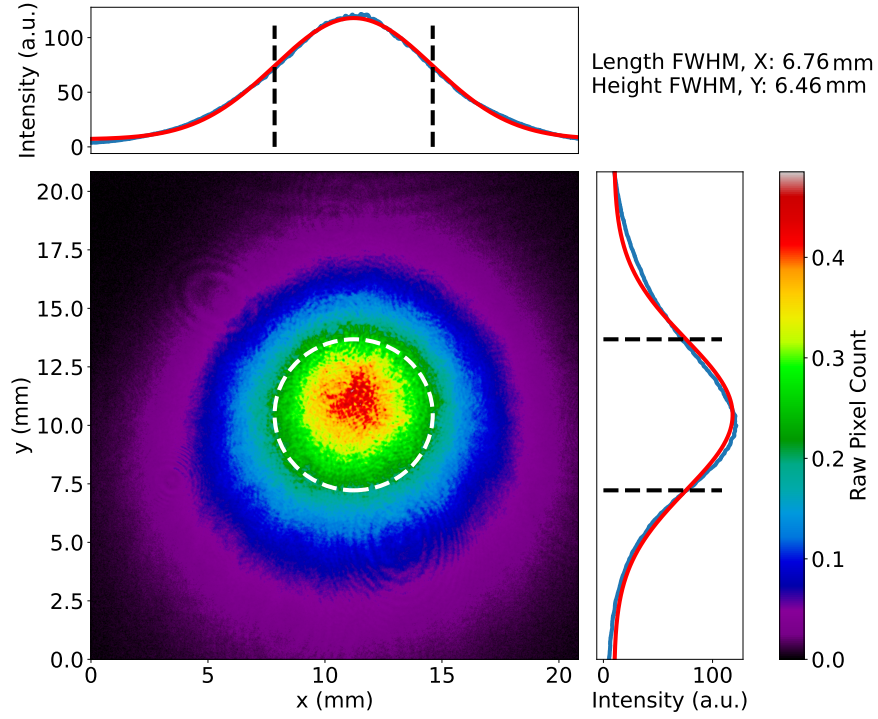


Figure 5.3: Coherent Legend near-field beam profile. Near-Field intensity distribution captured by CCD 1. The centroid is determined, and the white ellipse marks the $D4\sigma$ region. The horizontal and vertical line-outs are plotted on the top and side (light blue), and the FWHM size is determined by fitting a Gaussian function (red).

= 5 cm convex lens. From examining Figure 5.3, it is clear the pulse's intensity distribution is Gaussian along both axes.

The energy and energy stability were measured initially by a pre-calibrated Gentec QE95LP-H-MB-QED-D0 laser energy detector. This then allowed us to establish CCD 1 as an online laser energy diagnostic by quantifying the no. of CCD counts detected as a function of the laser energy. A mean background image for CCD 1 was recorded and subtracted from each laser light image to remove any noise or background light, and all images were cropped in post-processing so that only the portion of the CCD chip which is struck by the laser light is analysed. A test was performed to confirm that CCD 1 exhibited a linear response, in which the laser energy was varied by rotating the $\lambda/2$ waveplate to angles that were

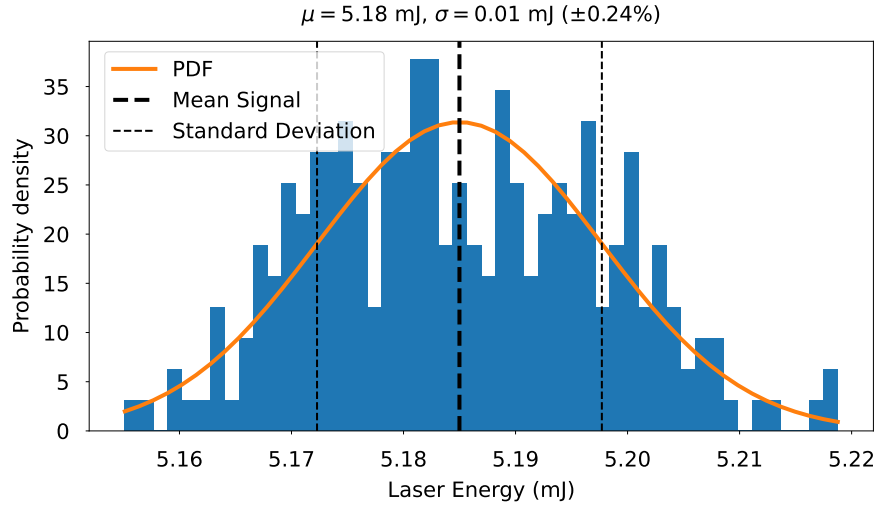


Figure 5.4: Laser energy stability. Histogram of the laser energy recorded by CCD 1 over 500 data points. μ and the central dashed line represent the mean value, where as σ and the exterior dashed lines denote the standard deviation. The orange line is a probability distribution function (PDF) fitted to the data.

demonstrated to correspond to specific laser energy values in the beamline when the laser energy detector was in place, and studying the behaviour of the summed intensity value recorded by CCD 1. Figure 5.4 shows a histogram of the laser energy and its probability density recorded over 500 separate data points which strongly resembles a normal distribution with a sub 0.5% standard deviation.

12 centimetres after passing through the chamber window, the laser pulse was focused by a $f = 15$ cm achromatic doublet lens to a focal spot size d_{FWHM} of approximately $7.23 \mu\text{m}$ in the horizontal direction and $5.89 \mu\text{m}$ in the vertical direction along a Rayleigh length of 215 mm, as shown in Figure 5.5. This profile was captured by attenuating the beam using a series of ND filters, placing a 10x Edmund Optics infinity corrected microscope objective at the focal position and then directing the light onto CCD 1, which was temporarily mounted inside the vacuum chamber, to form an image in the focal plane. The focus appeared to have a slight tail in the horizontal direction which was the result of an etalon

effect [246] creating multiple reflections of the laser pulse when passing through the $\lambda/2$ waveplate. This resulted in a small loss to the focal spot's overall encircled energy, which when considered alongside additional losses by the optics in the beamline and the spread in pulse duration due to dispersion effects, we estimated our peak intensity on target to be $I_0 \approx 6 \times 10^{16} \text{ W/cm}^2$ for our nominal laser parameters where the laser wavelength, $\lambda_L = 800 \text{ nm}$, laser energy, $W_L = 5.16 \text{ mJ}$ and laser pulse duration, $t_{FWHM} = 220 \text{ fs}$. This was selected as the reference case due to its capability to produce substantial afterglow across the entire gas density range, at an intensity value comparative to optical plasma torch injector lasers [117].

After generating the plasma filament, the laser is directed out of the chamber to intersect with a Polytetrafluoroethylene (PTFE) screen. PTFE is resistant to high temperatures and illuminates upon interacting with the laser, which allowed us to study the pulse profile and intensity post-interaction using CCD 4 and compare it to the image produced on CCD 1 for an identical data point. The laser intensity profile on the PTFE screen for our nominal laser parameters when the chamber is placed in air is shown in Figure 5.6. The slant in the profile from left to right is a result of the angle at which CCD 4 is observing the PTFE screen.

5.1.2 Beam transport

In order to gain control over the laser energy being supplied to the vacuum chamber, an attenuator was constructed using a $\lambda/2$ half waveplate and two polarising beam splitters. Since the beam splitters are designed to reflect 's' polarized and transmit 'p' polarized light, the light that exits this portion of the beamline could be adjusted by varying the polarization of the light incident on the beam splitter plates. This was achieved by placing the $\lambda/2$ half waveplate in front of the beam splitters and changing its angle of orientation, such that 0 degrees corresponded to fully 'p' (horizontal) polarized light and 45 degrees

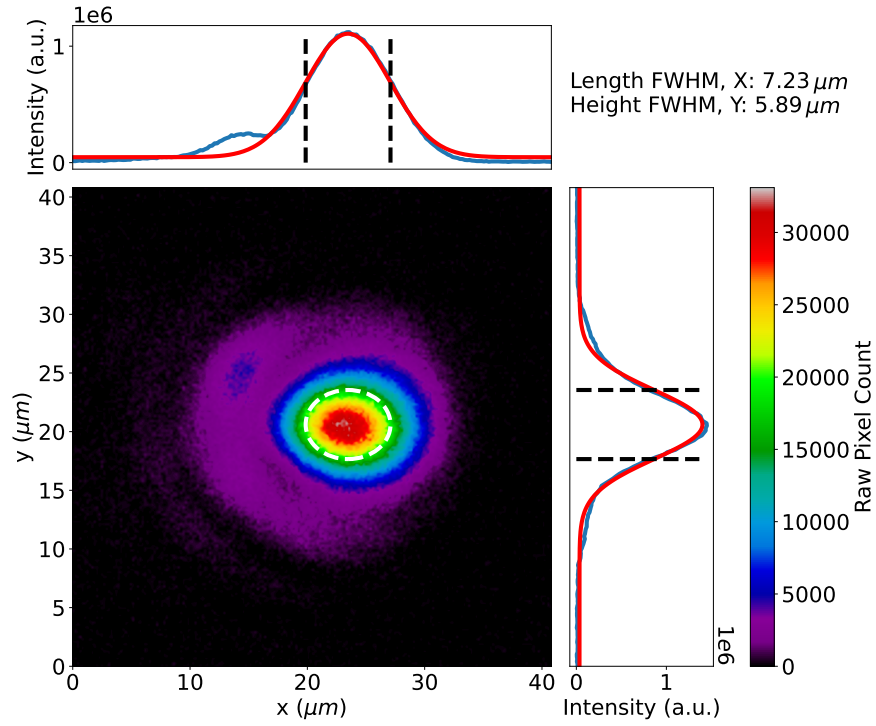


Figure 5.5: Coherent Legend far-field beam profile. Far-Field intensity distribution captured by a 10x Edmund Optics infinity corrected microscope objective. The centroid is determined, and the white ellipse marks the $D4\sigma$ region. The horizontal and vertical line-outs are plotted on the top and side (light blue), and the FWHM size is determined by fitting a Gaussian function (red). The horizontal aberration is the result of multiple reflections induced by an etalon effect in the $\lambda/2$ waveplate.

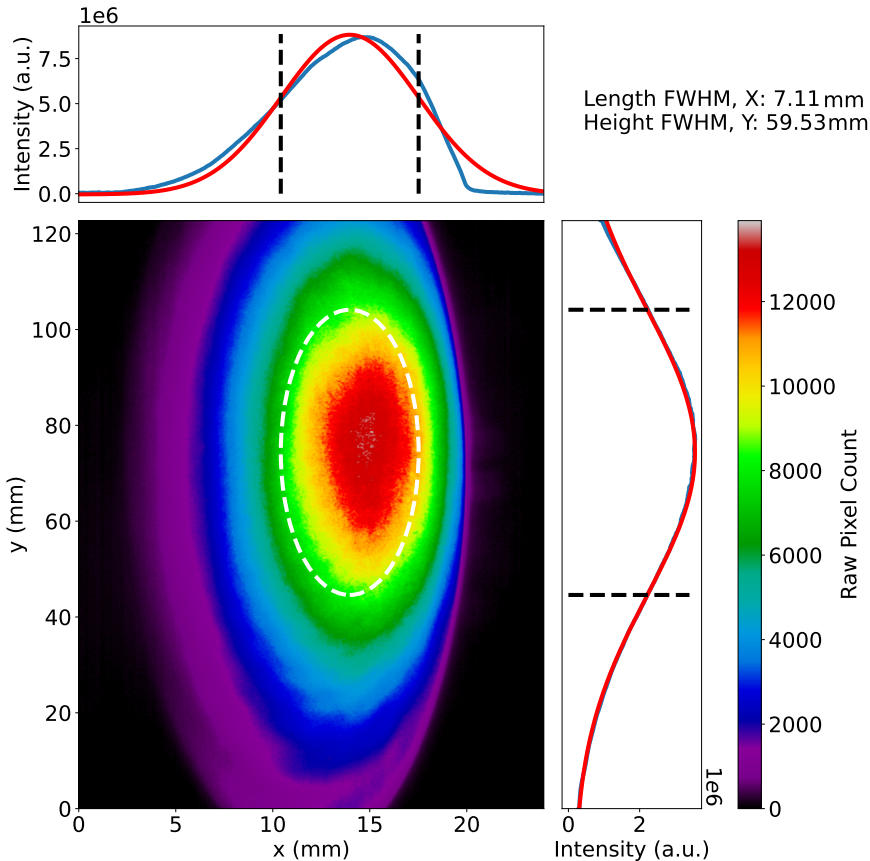


Figure 5.6: Coherent Legend PTFE screen profile. Laser intensity distribution on PTFE screen captured by CCD 4. The centroid is determined, and the white ellipse marks the $D4\sigma$ region. The horizontal and vertical line-outs are plotted on the top and side (light blue), and the FWHM size is determined by fitting a Gaussian function (red).

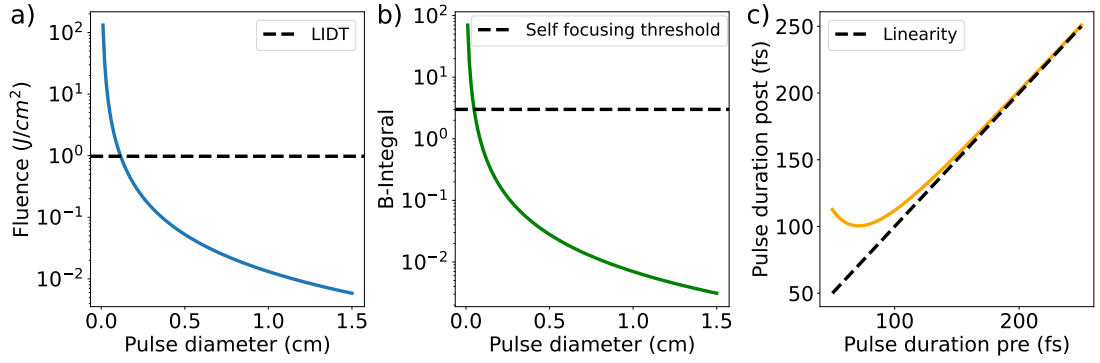


Figure 5.7: LIDT, B-integral and pulse lengthening clarification. Calculations to identify a set of laser pulse parameters that a) do not exceed the LIDT, b) undergo B-integral related self focusing or c) experience pulse lengthening as they traverse the 6.5 cm diameter CaF_2 vacuum chamber window port. The non-linear refractive index n_2 of CaF_2 is $1.26 \times 10^{-18} \text{ cm}^2/\text{W}$.

equated to fully ‘s’ (vertical) polarized light.

A series of mirrors were then used to transport the laser pulse along the beam-line towards the vacuum chamber. Any of the mirrors that interacted directly with the laser (M1-7) were ultrafast high-energy laser mirrors with laser-induced damage thresholds (LIDT) of $> 0.75 \text{ J}/\text{cm}^2$ at 800 nm with 150 fs pulses. The beam then passed through a 6.5 cm diameter calcium fluoride (CaF_2) window, coated to be anti-reflectant (AR) at 800 nm, to enter into the vacuum chamber. It was important to consider the potential for damage being inflicted on the window, along with any alterations to the laser that can be induced by the passage of a short, intense laser pulse through optical glass [247–249]. Graphical representations of calculations relating to the pulse’s intensity, duration, fluence $F = 2W_L/\pi r^2$ and B-Integral $B = \frac{2\pi}{\lambda_L} n_2 I_0$, where n_2 is the non-linear refractive index, are expressed in Figure 5.7, which helped to identify thresholds by which the aforementioned effects would either begin to occur or start to be prominently felt. Assuming the nominal beam diameter of 6.76 x 6.46 mm FWHM, the resulting fluence and B-Integral values sit comfortably below the LIDT and self focusing threshold. However for our designated window material and thickness,

it was determined that the pulse duration would have to be set at $t_{FWHM} \approx 220$ fs (see Figure 5.2 a)) in order to achieve < 0.5 % difference between the pulse duration pre and post window port and be deemed sufficiently resistant against pulse lengthening effects.

5.1.3 Experimental chamber

Helium was selected as the neutral gas medium of choice for this study due to its widespread use in LPA experiments over the last decade. The nominal neutral gas density was chosen to be $n = 5 \times 10^{17} \text{ cm}^{-3}$, as in combination with our reference laser parameters, this created a sufficiently bright afterglow signal that remained unaffected by non-linear effects present at higher gas densities. When operating in ‘flooded chamber’ mode, helium was released into the chamber using a needle valve, whilst monitoring the pressure via an Edwards Active Strain Gauge, which has ± 0.2 % accuracy across the full scale and was used to measure the pressure value during the parameter scans, and Wide Range Gauge tandem connected to an Edwards TIC Instrument Controller, and allowed to disperse homogeneously within the entire vacuum chamber. Therefore, once inside the chamber, the laser would propagate through helium gas at all times. The corresponding gas density value for each pressure was calculated using the ideal gas law [250], where the temperature inside the chamber was assumed to be identical to the kHz lab temperature as measured by a thermohygrometer. Although the majority of LPA experiments now employ a gas jet configuration, operating in ‘flooded chamber’ mode in the first instance not only provided a well-defined constant gas density reservoir which was advantageous for minimising any fluctuations to the plasma afterglow when trying to map out the signal’s dependency on experimental parameters, but also the ability to access densities that are beyond the standard range of gas jets.

Alternatively, gas jet mode consisted of helium being released from the jet

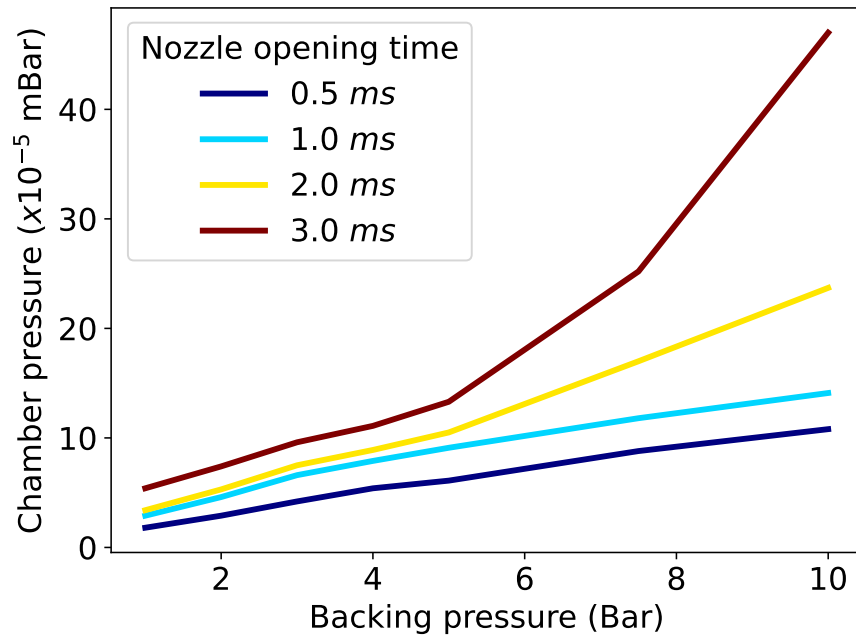


Figure 5.8: Pumping system performance. Analysis of the pumping system performance by monitoring the vacuum chamber pressure for a select combination of backing pressures and valve opening times.

into a small volume above its opening, through which the laser would traverse, for a defined period of time at a defined rate. This was controlled externally using a Parker IOTA ONE pulse driver. In order to synchronize the arrival of our laser pulse and the opening of the gas jet, a Stanford Research Systems DG645 digital delay generator was utilised. A signal from our laser system was fed in as the external trigger source allowing for several precisely defined pulses to be created that had a specific length (in time) and delay with respect to the laser pulse. This not only enabled successful gas jet operation, but was essential to ensure each diagnostic device was recording in unison to capture the same afterglow signal. In this mode, the gas density was acquired by simulating the flow through the nozzle for each respective pressure using the modeling software Ansys Fluent [251] (see Section 5.4).

Both operational modes required the vacuum chamber to first be placed under

vacuum before shooting. The pumping system adopted was a combination of one primary dry pump (Edwards EDS300) and one turbo pump (Edwards STP Maglev) with suction from 300 l/s to 4,500 l/s. With just the dry pump running, a vacuum level of 1×10^{-3} mbar was achievable. If data collection was to be performed in flooded chamber mode, the vacuum chamber would be sealed off from the pump and placed in a steady state before allowing the helium to flow in. Trace amounts of air was able to filter into the vacuum chamber during each scan due to minute leaks at the chamber doors and/or window ports, which can be visualised in the integrated afterglow spectrum seen in Figure 5.11 b). However, the percentage of air inside the vacuum chamber is dwarfed when compared to helium, meaning it is effectively inconsequential to the laser-plasma interaction and ensuing afterglow signal. If a gas jet scan was to be undertaken, the vacuum chamber would not be sealed off and the turbo pump would be switched on, allowing a vacuum level of 2×10^{-6} mBar to be achieved. Before carrying out any gas jet scans, a circular 3 mm SourceLAB SL-NOZ I nozzle, designed specifically for electron acceleration, was attached to the valve and the jet was pulsed for the range of different backing pressures - 1 to 10 Bar, and opening times - 0.5 to 3.0 ms, due to be explored in the experiment, whilst monitoring the vacuum chamber pressure on the Turbo and Instrument Controller (TIC). This was done to establish whether the pumping system could comfortably accommodate the gas load required to perform the experiment whilst still keeping the vacuum chamber pressure at an ‘acceptable’ level, which was deemed to be below 1×10^{-3} mBar. The chamber pressure was recorded approximately 30 seconds after the pulsing began to allow time for the value to stabilise. Both the gas jet repetition rate and CCD acquisition rate were set to 1 Hz, which equates to a period of $T = 1/f = 1$ s, during the experimental campaign to mitigate against a rise in the chamber pressure and allow the camera software sufficient time to process and save each recorded image between data points. With the laser still operating

at kHz repetition rate, this meant one every one thousand generated afterglow signals was captured. Figure 5.8 shows the test results for the chosen combination of backing pressures and opening time cases. It is clear the chamber pressure remains below the designated threshold for the range of gas jet parameters to be explored during the subsequent parameter scans.

A final check performed before beginning any parameter scans was to establish the optical ‘thickness’ of the helium gas inside the vacuum chamber when the experiment is placed in ‘flooded chamber’ mode. This was important to clarify so as to see whether reabsorption of the afterglow light as it traversed the vacuum chamber had to be taken into consideration or whether it is negligible and can be ignored. The Thorlabs’ SLS201/M lightsource was positioned outside the window port on the left-hand side of the chamber and shone through the vacuum chamber towards CCD 2, which was briefly re-oriented to face the light source. 50 images were recorded of the light source at different chamber settings - vacuum, 10^{16} , 10^{17} , 10^{18} & 10^{19} cm^{-3} , with the mean summed intensity value for each calculated and compared. There was found to be only 0.4% difference between the highest and lowest intensity values. Light intensity decreased between vacuum and 10^{17} cm^{-3} but then increased between 10^{17} & 10^{19} cm^{-3} . This coupled with the small variance suggests that these shifts are just fluctuations and not a trend associated with reabsorption of the light at higher densities.

5.1.4 Laser-gas interaction

Along the direction of focus as the pulse intensity increased, the laser began to optical-field ionize the helium gas over a period of hundreds of femtoseconds and generate a narrow plasma filament. To better understand the laser-gas interaction, high fidelity FBPIC simulations were performed to replicate the experiment. The simulation box size was set at $345.5 \times 924.8 \text{ } \mu\text{m}$, with the cell size in z , $D_z = 5 \times 10^{-8} \text{ } \mu\text{m}$, the cell size in r , $D_r = 4 \times 10^{-7} \text{ } \mu\text{m}$ and the number of modes used,

$N_m = 2$. A moving window setup was employed with the window's velocity set to the speed of light, c . The number of particles per cell along each axis were as follows - $p_{nz} = 1$, $p_{nr} = 2$ & $p_{nt} = 4$. Stationary helium atoms were added as the particle species, with each macroparticle capable of being ionised via the ADK ionisation model (see Section 2.3.1). The laser pulse was set to propagate for 4 mm, with the focus set at the midway point, 2 mm.

In the first instance, the simulation was performed for our nominal case (see Sections 5.1.1 and 5.1.3) with the resulting plasma filament and spectra of ionized electrons that encompass the filament displayed in Figure 5.9. The filament visibly pinches down at its centre mimicking the spatial profile of the focused Gaussian laser pulse creating it, and extends well outwith the laser's Rayleigh length in z , due to the large amount of energy present in each pulse, whilst being spatially confined to tens of microns in y . As seen from the spectrum, the interaction gives rise to a 'hot' plasma with liberated electrons from both ionization levels and particles that possess initial temperatures of up to 6000 eV, with the higher laser intensity required to liberate He 2+ electrons being translated into higher electron temperatures for this species. Notably, a large quantity of the electrons are produced with kinetic energies of ≈ 66 eV, where the cross section for electrons scattering off helium gas is peaked, meaning that a substantial amount of secondary impact ionization events are likely to take place, generating more electrons. However, FBPIC does not contain an impact ionisation module at present and was therefore unable to simulate this secondary ionisation. Similarly, separate modelling codes are required to replicate the plasma expansion and light emission phases of the plasma's lifecycle [252, 253].

Subsequently, thermalisation of the plasma electrons occurs, tens of picoseconds after their genesis [254, 255], followed by a phase of expansion in which the plasma moves outwards over the course of a few nanoseconds [43, 256, 257]. Then comes a period, tens of microseconds after the laser has passed, in which

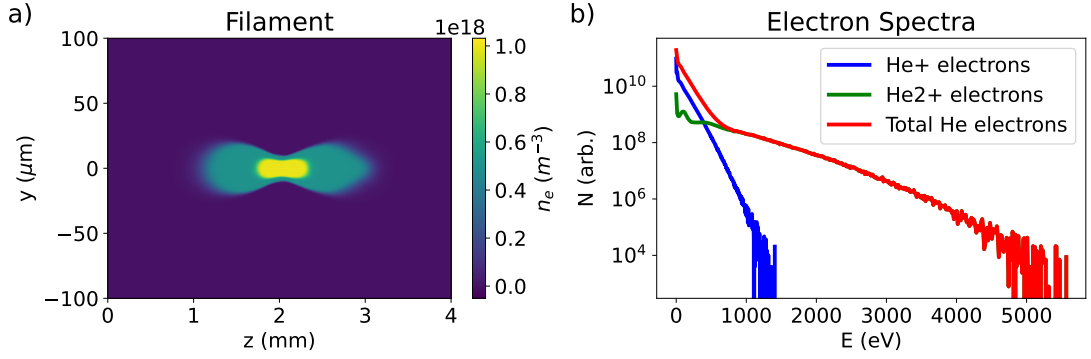


Figure 5.9: Particle-in-cell simulated plasma filament and electron spectra. FBPIC simulated initial a) plasma filament snapshot, generated by a laser pulse polarised in the x direction, propagating from left to right where n_e represents the electron density, and b) its associated electron energy spectrum, where only He 2+ electrons have energies > 1500 eV, hence the merging of the green and red lines at this energy.

more gradual thermalisation and diffusion between the plasma and neutral gas takes place. During this time, electrons either recombine with the plasma ions or collide with the neutral atoms to excite them, with the former leading to the complete neutralisation of the plasma after some hundreds of microseconds [48, 258, 259]. Both of these processes emit radiation which together with potential bremsstrahlung radiation, depending on the chosen plasma conditions, comprise the plasma afterglow signal. This emission contributes to plasma cooling and contains valuable information about the initial laser-gas interaction and the plasma's physical properties.

5.1.5 Afterglow diagnostics

Various diagnostics were used to study the afterglow signal during the experiment. CCDs 2 and 3 were fitted with Edmund Optics VariMag Telecentric Lenses, that were adjusted accordingly to bring the afterglow into focus, and oriented with the aid of protected silver mirrors, with high reflectivity in the visible spectrum, to observe the afterglow from the side (along the x axis), in the case of CCD 2, and

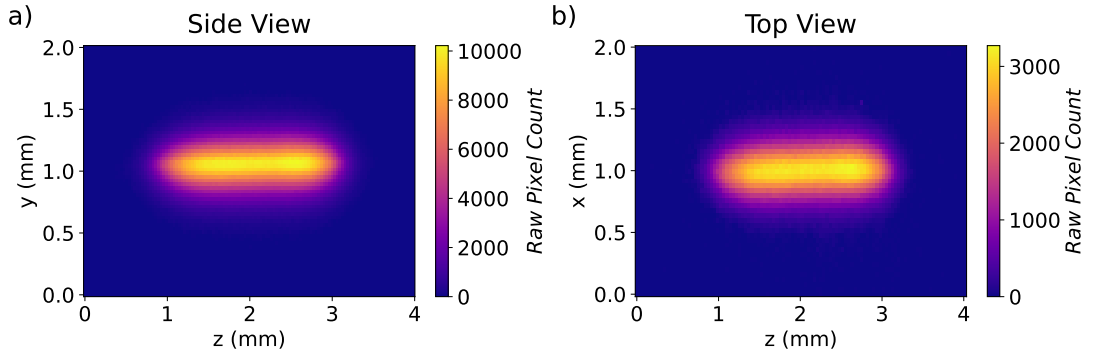


Figure 5.10: Plasma afterglow - side and top view. Experimentally recorded images of the mean plasma afterglow signal from 50 data points, as seen from both the a) side, z-y plane, and b) top, z-x plane, view. The laser is polarised in the z-x plane.

from the top (along the y axis), in the case of CCD 3. Images of the afterglow signal as seen from these two cameras are shown in Figure 5.10.

Additionally, an Andor Shamrock 303i Spectrograph was installed outside the window port on the left-hand side of the chamber which enabled the afterglow light to be spatially spectrally resolved. Connected to the Shamrock was an Andor iDus camera which boasted a high dynamic range, high quantum efficiency and a temperature cooling system that allowed for dark current to be effectively suppressed. To calibrate the wavelength of the spectrometer, an Ocean Optics' HG-1 Mercury-Argon lamp was used in combination with a third-order linear regression calculator to assign a wavelength value to each camera pixel. The resolution was calculated to be 0.560 nm/px over a wavelength range of 390 - 760 nm. As with the CCD's, the spatial calibration for the Andor iDus was determined by inserting a measuring tape into the imaging system and observing the number of pixels that the object encompassed, which gave rise to a resolution of 0.038 mm/px. The quantum 'yield', or efficiency, of the spectrometer at each wavelength was also calculated using a Thorlabs' SLS201/M stabilized Tungsten-Halogen light source, which provides a constant-intensity blackbody radiation spectrum from 360 to 2600 nm. This calibration was then applied to each recorded

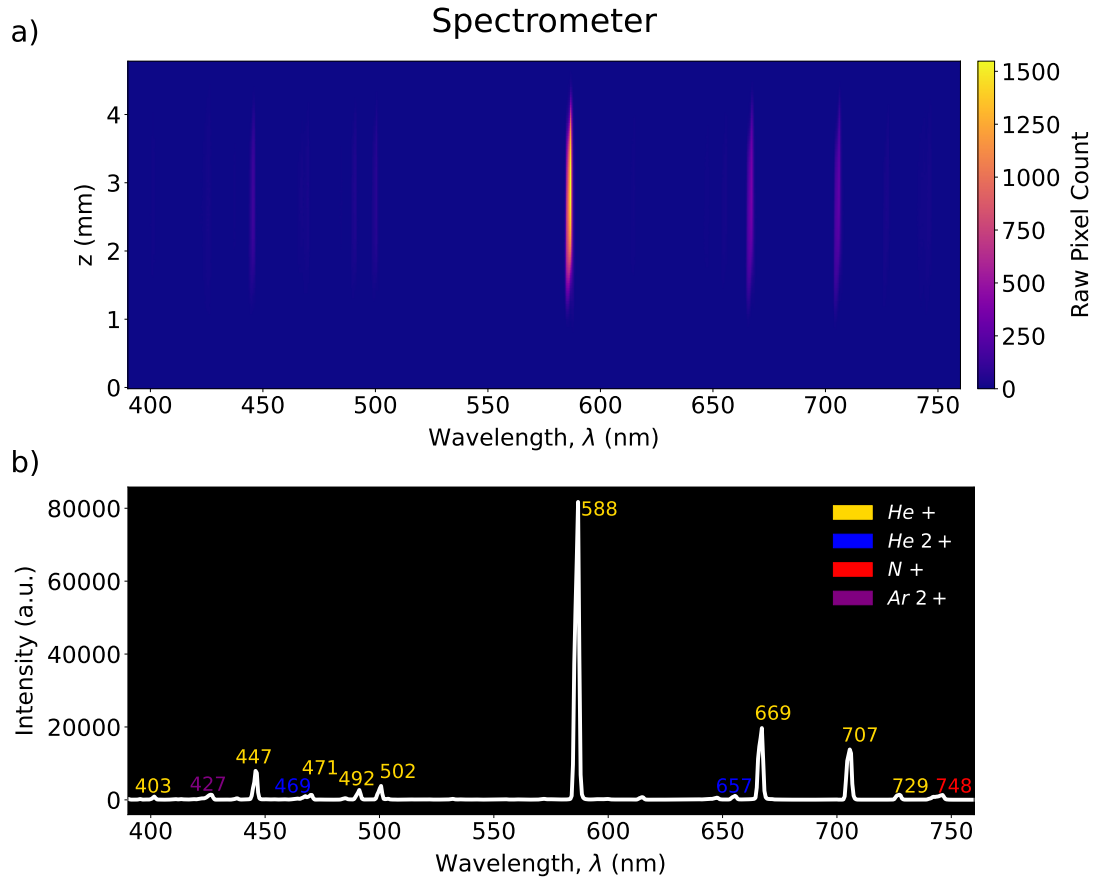


Figure 5.11: Spatial and integrated afterglow spectrum. Experimentally recorded images of the mean plasma afterglow light spectrum from 50 data points a) spatially resolved along the laser propagation axis, z , and b) integrated with each emission line assigned to the responsible gas species and ionization level.

spectrum during post processing so as to be able to more accurately compare the intensity shifts of individual spectral lines during a parameter scan.

The afterglow light travelling out of the chamber window was imaged using an $f = 20$ cm achromatic doublet lens onto the entrance slit. A dove prism, that enabled for the orientation of the incoming afterglow light to be both rotated and flipped, was positioned just in front of the Shamrock and mounted on a mechanical stage, giving the freedom to either insert or remove the prism depending on the requirements of a specific scan. Figure 5.11 demonstrates a typical Shamrock

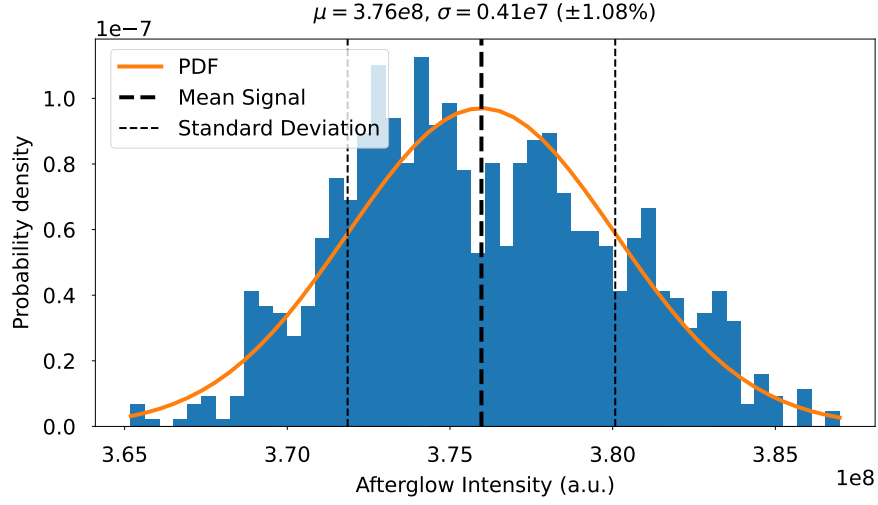


Figure 5.12: Afterglow intensity stability. Histogram of the afterglow intensity recorded by CCD 2 over 500 data points. μ and the central dashed line represent the mean value, where as σ and the exterior dashed lines denote the standard deviation. The orange line is a probability distribution function (PDF) fitted to the data.

output image along the laser propagation axis z , alongside a spatially integrated spectrum of the light collected for our nominal case - $\lambda_L = 800$ nm, $W_L = 5.16$ mJ, $t_{FWHM} = 220$ fs and $n = 5 \times 10^{17}$ cm $^{-3}$, where each line that extends above the background noise has been attributed to its parent process. The majority of these lines correspond to strong lines within the helium emission spectrum, with the few exceptions being the 427 nm Ar 2+ and 748 nm N+ lines that originate from the trace amounts of air in the vacuum chamber as previously discussed. Each one of our afterglow diagnostics was equipped with a 750 nm OD = 4.0 shortpass filter to eliminate any light from the laser's spectrum reaching the cameras.

In a similar style to the input laser energy shown in Figure 5.4, Figure 5.12 is a histogram of the afterglow signal's intensity and its probability density over 500 consecutive data points. As with CCD 1, mean background images for CCD 2, CCD 3 and the Andor iDus were recorded and subtracted from each respective afterglow image to remove any noise or background light, and all images were

cropped in post-processing so that only the portion of the CCD chip which is struck by the afterglow light is analysed. An additional function is used in post processing to find and remove any hot or dead pixels, with the user defining the number of standard deviations used as a cut off limit. This plot showcases the $\approx 1\%$ standard deviation achieved, which is an order of magnitude better stability than was realised during the E210 laser only runs at SLAC FACET I [116]. Notably, the afterglow intensity jitter is 4.5 times greater than laser energy jitter, which is to be expected as many other experimental parameters contribute to the afterglow signal such as the focal spot, pulse duration etc. which all fluctuate by varying degrees on a shot to shot basis.

5.2 Experimental results - Flooded chamber mode

In order to establish afterglow dependency and gauge sensitivity to variance of experimental parameters, scans were performed in which the chosen quantity was varied by known increments, with the resulting afterglow signal recorded at each step. These specific parameters: laser pointing angle, laser energy and gas density, were selected as they can each significantly impact the laser-gas interaction and it was possible to discern the degree to which each quantity was fluctuating on a shot-to-shot basis using the equipment available within the lab.

During the series of parameter scans, observations were made of the afterglow light that warrant some discussion and further analysis up front. What was immediately apparent when comparing the afterglow signal (see Figure 5.10) with the initial laser generated plasma filament displayed in Figure 5.9, is that the emitted light extends out further radially, as it encompasses the various secondary ionization processes and expansion phase of the plasma's life cycle described in Section 5.1.4.

Furthermore, when comparing the top and side view images of the afterglow

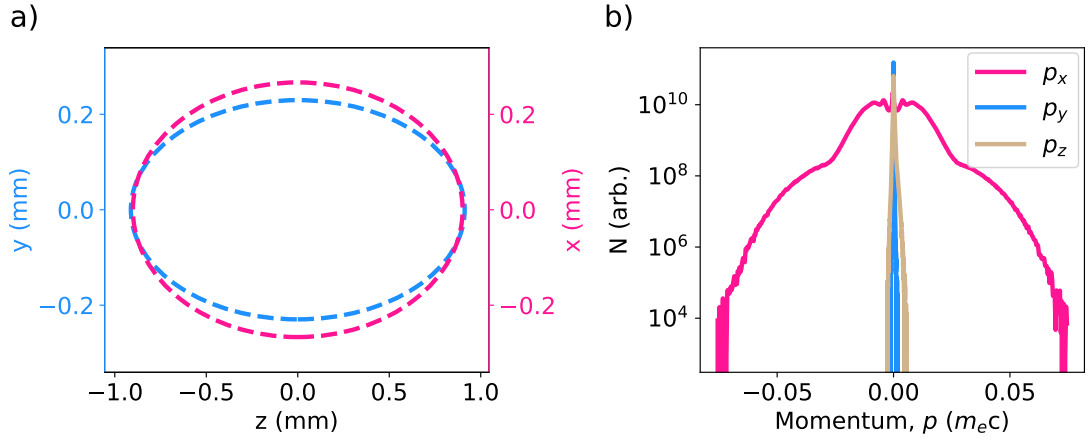


Figure 5.13: Asymmetric afterglow signal. a) The afterglow light was identified to extend out further along the z-x axis (top view) than the z-y axis (side view), as seen here for the nominal parameters case, mimicking the expected asymmetric profile of a plasma generated by laser pulse linearly polarized in the x direction. b) The initial momentum of the FBPIC simulated laser ionized electrons in each direction is displayed, highlighting that the majority of their kinetic energy is directed along the x plane.

recorded by CCD's 2 & 3 (see Figure 5.10), the light was seen to consistently extend out further along the z-x axis when compared with the z-y axis, as seen for the nominal parameters case in Figure 5.13 a). This mimics the expected asymmetric profile of a plasma generated by a laser pulse polarized in the x direction [169] as electrons born earlier will travel further to the outside of the plasma volume than electrons born later, and in this case the first born electrons generated by laser optical field ionization (OFI) possess the majority of their momentum in the x direction as seen in Fig 5.13 b), hence the thermal expansion along this axis. This trend was consistent for all parameter cases. It is notable these asymmetries in the initial plasma profile are still reflected in the afterglow signal despite the separation of several in-between plasma lifecycle phases. This highlights the importance of observing the afterglow from multiple viewpoints in future experimental setups in order to fully characterise the signal's evolution.

Another interesting detail to arise from the experimental scans was that by em-

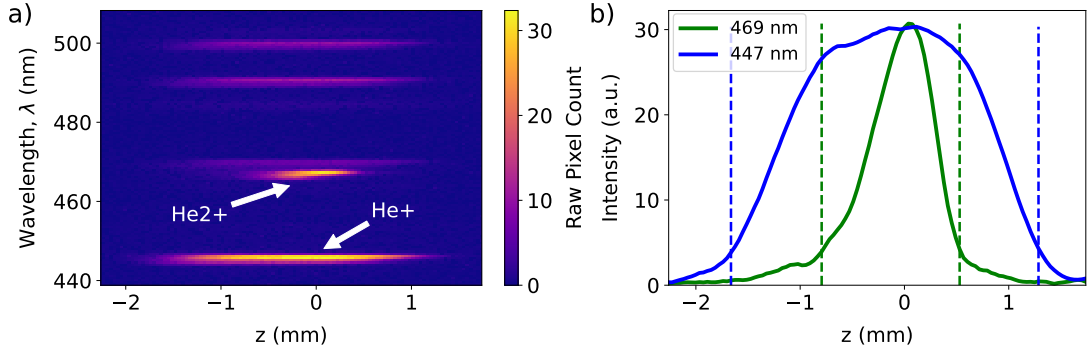


Figure 5.14: Spatially mapped ionization regions. A combination of the shamrock spectrometer and dove prism allows for the afterglow signal to be broken down into the portion of light emitted by each ionization level and how far along the propagation, z , this light extends. Subplot a) demonstrates this using the 447 nm He^+ and 469 nm He^{2+} emission lines and subplot b) showcases a lineout of these signals (solid lines) with the respective dashed lines representing the $1/e^2$ length.

ploying the shamrock spectrometer, ionization regions of multi-level gas species, like Helium, can be mapped out by spatially spectrally resolving the afterglow light. This is shown in Figure 5.14 where the mean afterglow image recorded by the Shamrock spectrometer from 50 data points for our nominal case - $\lambda_L = 800$ nm, $W_L = 5.16$ mJ, $t_{FWHM} = 220$ fs and $n = 5 \times 10^{17} \text{ cm}^{-3}$ is displayed. This opens up the potential to obtain a longitudinally resolved measurement of the laser intensity along the focus for specific cases. For instance, low laser intensity and low gas density that will generate a ‘cold’ plasma filament would lead to minimal impact ionization and plasma expansion and therefore the afterglow’s longitudinal profile would be more representative of the optical field ionized plasma.

5.2.1 Laser pointing angle variation

The first parameter of choice was the beam pointing of the laser pulse, which was determined in the first case by venting and then opening the vacuum chamber,

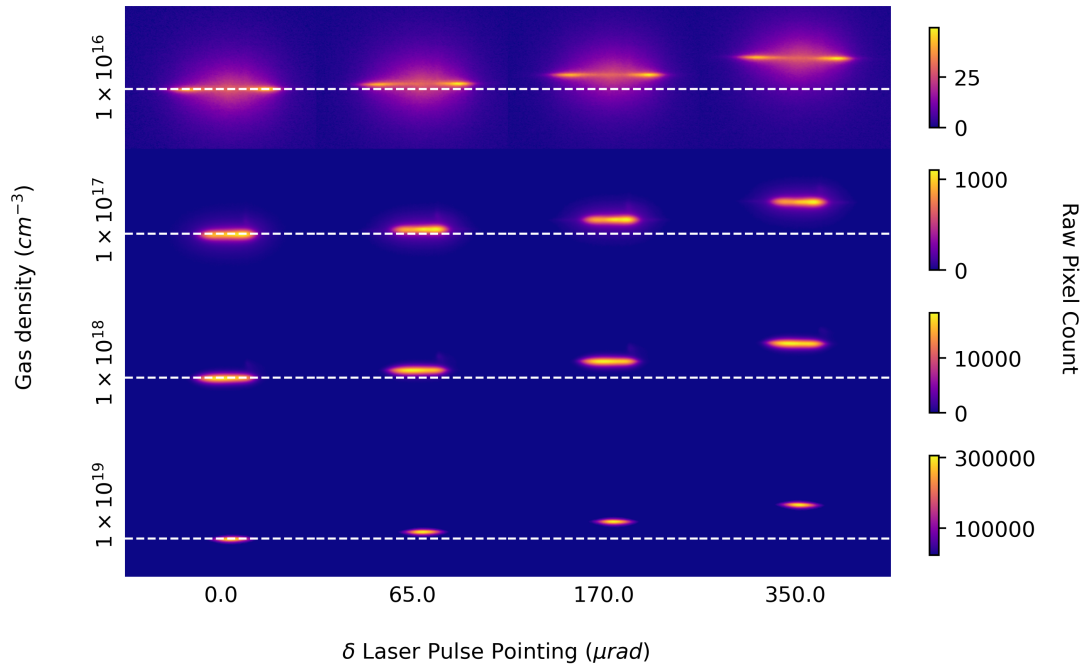


Figure 5.15: Afterglow Images - Laser pointing angle variation. Images of the mean afterglow signal recorded by SideView CCD 2 from 50 data points for 4 different laser pulse pointing angles - 0, 65, 170 and 350 μrad , and 4 gas density cases - 10^{16} , 10^{17} , 10^{18} & 10^{19} cm^{-3} . The white dashed line represents the nominal y position.

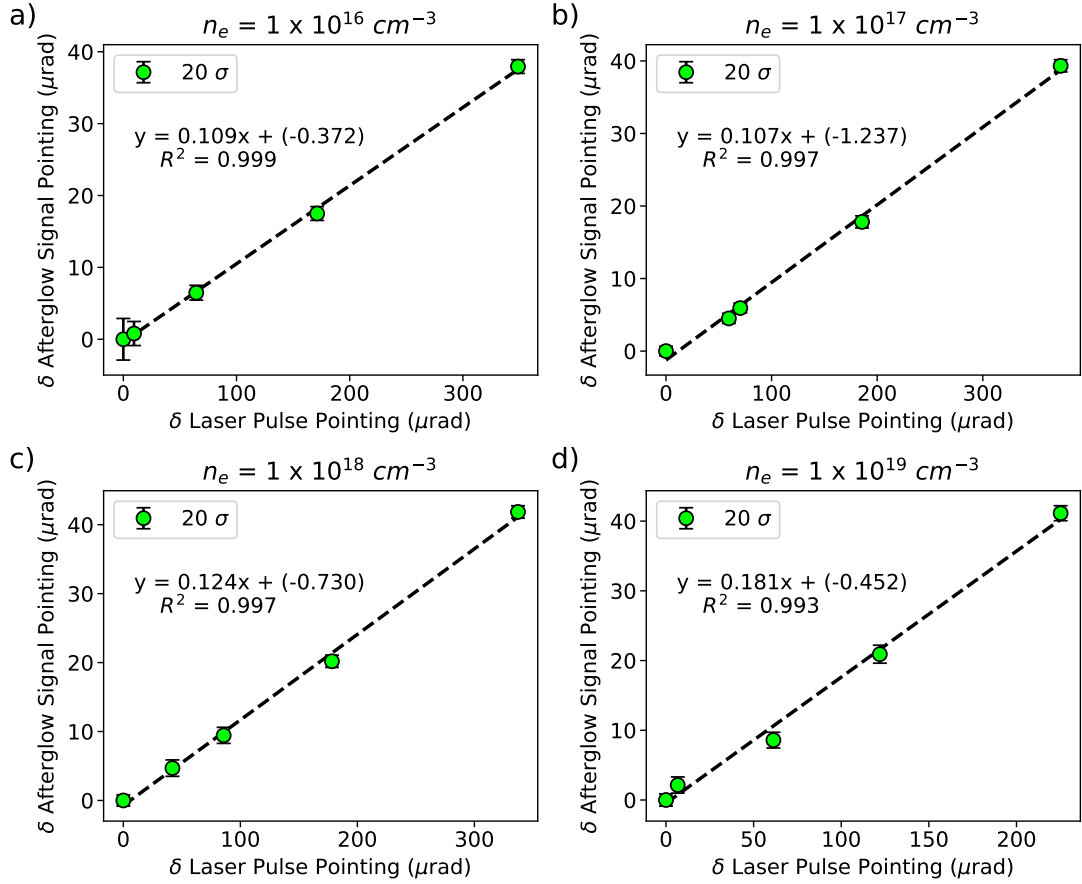


Figure 5.16: Laser spatial movement scan. Laser pulse's pointing is varied and the corresponding afterglow signal's pointing is recorded (green markers) at 4 different gas densities - $n =$ a) 10^{16} b) 10^{17} c) 10^{18} & d) 10^{19} cm^{-3} . Each marker represents the mean afterglow signal pointing for 50 data points. A linear regression model is applied to the data (dashed line) with the line of best fit equation and R^2 value displayed to quantify the proximity of our data to the fitted regression line. The error bars encompass 20 times the standard deviation, σ .

adjusting the screws on the mirror mount containing M6, the last mirror before the laser pulse enters the vacuum chamber, in order to change the distance between the front and back plate on the mount by $\approx 50, 100, 250$ and 500 microns which consequently tilted the beam upwards, and then examining the beam's change in position on the PTFE screen. This coupled with measuring the path length from M6 to the PTFE screen, allowed for the laser's pointing angle to be quantified. The path length from M6 to the laser focus was also measured so that the corresponding afterglow pointing angle could later be calculated using the same method.

The vacuum chamber was then closed and flooded with helium by following the procedure detailed in section 5.1.3, before adjusting the M6 mirror mount by the same amounts as laid out above and recording the corresponding afterglow signal. This scan was performed with the nominal laser parameters - $\lambda_L = 800$ nm, $W_L = 5.16$ mJ and $t_{FWHM} = 220$ fs, and at 4 different gas densities - $n = 10^{16}, 10^{17}, 10^{18}$ & 10^{19} cm $^{-3}$. The mean afterglow images recorded by SideView CCD 2 from 50 data points are displayed in Figure 5.15 for 4 different laser pulse pointing angles - $0, 65, 170$ and 350 μ rad, and all 4 gas density cases. As the pointing angle is increased, we can see the afterglow signal respond to this by shifting upwards from the nominal y position, represented by the white dashed line. Figure 5.16 details the relationship between the change, δ , in the laser pulse pointing and the change in the afterglow signal pointing for each of the four gas density cases. Clearly visible from these plots is the high degree of linearity between the laser pulse's positional movement and the afterglow's positional movement when the laser pointing is changed, confirming that the afterglow is in fact sensitive to spatial changes experienced by the laser. There is approximately a factor of 10 difference between the two quantities, caused by the magnitude of laser pulse pointing being lessened in the focal plane, where the afterglow signal originates. It is also worth drawing attention to how minimal

the standard deviation is for the afterglow pointing jitter measurements, with the corresponding error bars having to cover 20σ instead of 1σ so that they are visible on the plots. An improved camera resolution, where the afterglow signal covers more pixels on the CCD chip, would almost certainly lead to a larger error however it is nonetheless valid to emphasise the stability of this measurement.

The R^2 value of the trend lines, which is displayed on each plot and represents how well our data fits the applied linear regression model, decreases slightly as the gas density is increased, from 0.999 at 10^{16} cm^{-3} to 0.993 at 10^{19} cm^{-3} . This suggests the afterglow may be less sensitive to the laser's movement at higher densities where the laser-gas interaction became increasingly non-linear. However, this cannot be confidently inferred from the experimentally recorded data given the change in R^2 value is relatively small. Ultimately, confirmation of this linear relationship allows the afterglow signal to be employed as a high-precision diagnostic for measuring the laser pointing angle at the interaction point. This could be useful for laser-plasma-interactions where the incoming laser beam pointing changes the position where the laser interacts with a gas target, and may therefore change the trajectory of produced electron beams [188].

5.2.2 Laser energy variation

Next, we chose to explore the afterglow's response to changes in the laser's peak intensity, by way of varying the laser energy (see Equation 2.19) using the attenuator setup outlined in Section 5.1.2. As described in Section 2.3.1, tunneling ionization is highly sensitive to the laser pulse's electric field and peak intensity. Therefore, if the afterglow exhibits a linear response, it may allow one to exploit the signal as an indestructible laser intensity diagnostic, directly at the interaction point for constant gas densities. Initially the laser energy was adjusted by approximately 0.05, 0.1 & 0.2 mJ from the nominal case - $W_L = 5.16 \text{ mJ}$, which constituted 0.97, 1.94 & 3.88 % of the total energy of the beam, with the chamber

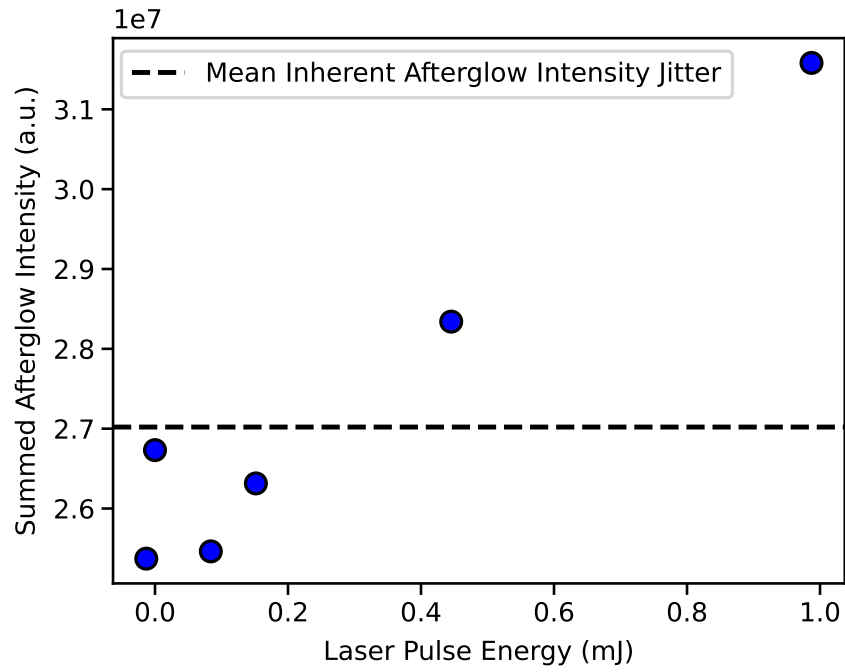


Figure 5.17: Small range laser energy scan. Laser pulse's energy is varied by fine amounts and the corresponding summed afterglow signal's intensity is recorded (blue markers). Each marker represents the mean summed afterglow intensity for 50 data points. The dashed line represents the mean inherent intensity jitter associated with the afterglow signal.

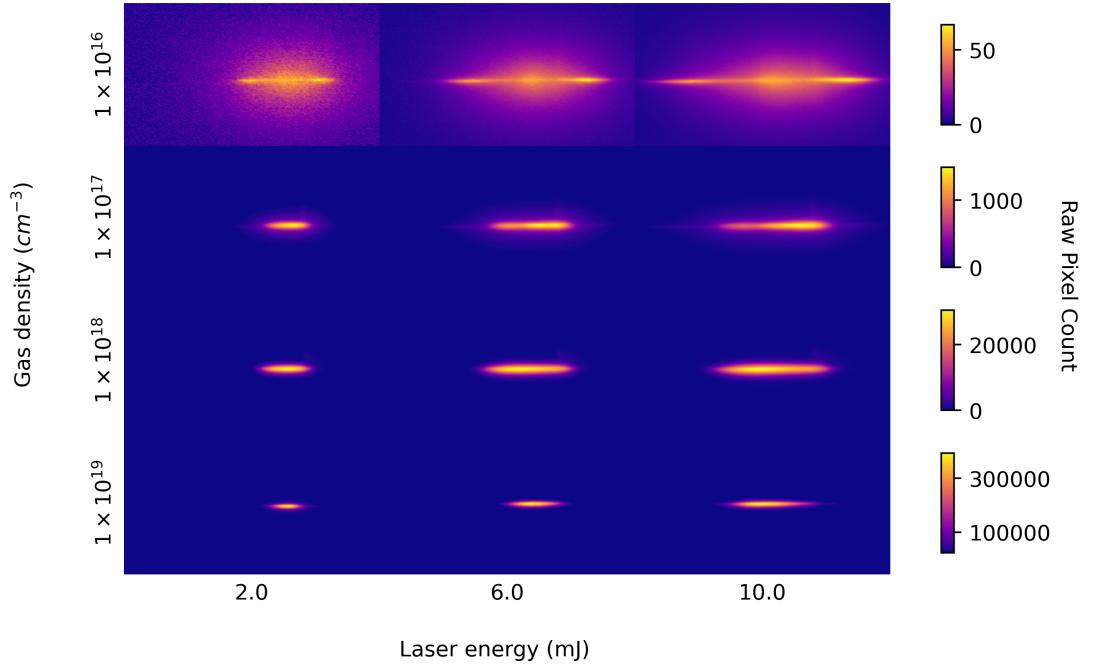


Figure 5.18: Afterglow Images - Laser energy variation. Images of the mean afterglow signal recorded by SideView CCD 2 from 50 data points for 3 different laser energies - 2, 6 and 10 mJ, and 4 gas density cases - 10^{16} , 10^{17} , 10^{18} & 10^{19} cm^{-3} .

flooded to the nominal gas density - $n = 5 \times 10^{17} \text{ cm}^{-3}$, and the corresponding afterglow signal being measured at each step. The laser pulse energy plotted against the summed afterglow signal intensity for these four cases can be seen in the bottom left corner of Figure 5.17.

There appears to be no obvious trend between these small increases in the laser energy and the afterglow intensity. This can however be explained by referring back to Figure 5.12, which displays the stability of the afterglow signal observed during these measurements. Despite minimizing the inherent afterglow intensity jitter here to just 1.08%, this can still impact on any potential measurements that don't bring about a change greater than this inherent jitter. Essentially, the change in the chosen parameter, which in our case here was the laser energy, must induce a change in the afterglow intensity that is above the inherent afterglow

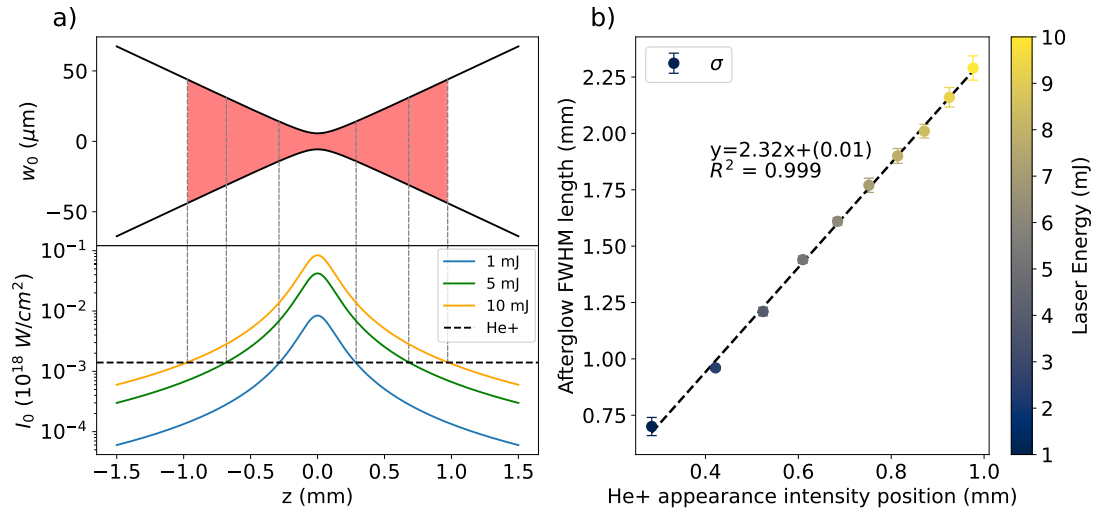


Figure 5.19: Laser energy - Afterglow FWHM length correlation. Explanation of the rate at which the afterglow FWHM length increases with the laser energy. The theoretical intensity profile of a Gaussian laser pulse along z is shown in a) for 3 laser energy cases – 1 mJ, 5 mJ & 10 mJ, alongside the corresponding ionization region as indicated by the vertical grey dashed lines travelling upwards from the z position at which the laser pulse exceeds the He+ appearance intensity threshold I_{app} (horizontal black dashed line). In b), the afterglow FWHM length for each laser energy case, calculated by taking the mean value across 50 experimental data points, can be seen to scale linearly with the He+ appearance intensity position. The data markers are colour coded in conjunction with colour bar on the right-hand side of the plot, and the error bars represent the standard deviation, σ .

intensity jitter threshold (dashed black line in Figure 5.17). If it doesn't, you will likely see no obvious correlation as the trend line is susceptible to these inherent fluctuations that can mask the intrinsic relationship between the afterglow and selected parameter. These inherent variations in the afterglow can be likened to the noise floor of any standard diagnostic device, for instance the noise level of a typical CCD, and will vary depending on the overall stability of your system. This same principle is also applicable to the afterglow's spatial movement discussed in 5.2.1, however it is not visible as the induced variations far exceed the inherent ones in the scan carried out.

As shown in the middle and top right portion of Fig 5.17, when the mean inherent afterglow intensity jitter is exceeded, in this case through increasing the laser energy by 0.5 & 1.0 mJ, the correlation between the laser energy and the afterglow intensity can start to be described by a linear function. This was then expanded upon by performing a follow up scan over a larger dynamic range where the laser energy was varied from $W_L = 1 - 10$ mJ, at 4 different gas densities - $n = 10^{16}, 10^{17}, 10^{18} \& 10^{19} \text{ cm}^{-3}$. The mean afterglow images recorded by SideView CCD 2 from 50 data points are displayed in Figure 5.18 for 3 different laser energies - 2, 6 and 10 mJ, and all 4 gas density cases.

From observing the afterglow light in these images, the length of the signal along z clearly increases as the laser energy is increased. This aligns with our assumption that since a laser pulse containing more energy is able to ionize further along the propagation axis, z , the afterglow light emanating from this plasma should reflect this and also increase in length. On further analysis, the afterglow FWHM length for each laser energy case was found to scale linearly with the z position at which the laser pulse exceeds the appearance intensity I_{app} for He+ as shown in Fig 5.19 b) for our nominal gas density - $n = 5 \times 10^{17} \text{ cm}^{-3}$ case. This confirms that the laser's Gaussian intensity profile along z , as seen in 5.19 a) for 3 laser energy cases - 1, 5 and 10 mJ, is reflected in both the plasma filament and subsequent afterglow signal.

Figure 5.20 details the relationship between laser pulse's energy and afterglow signal's intensity for each of the four gas density cases. It is clear there is a high degree of linearity between the two quantities, and thus it is possible to use the afterglow to accurately diagnose changes in the laser pulse's energy/peak intensity. This trend becomes marginally less linear in the high density regime, as was the case for the laser pointing angle scan outlined in Section 5.2.1. When operating with gases that have multiple ionization levels (He+ to He2+ in this instance), having an awareness of the appearance intensity, I_{app} , for each ion

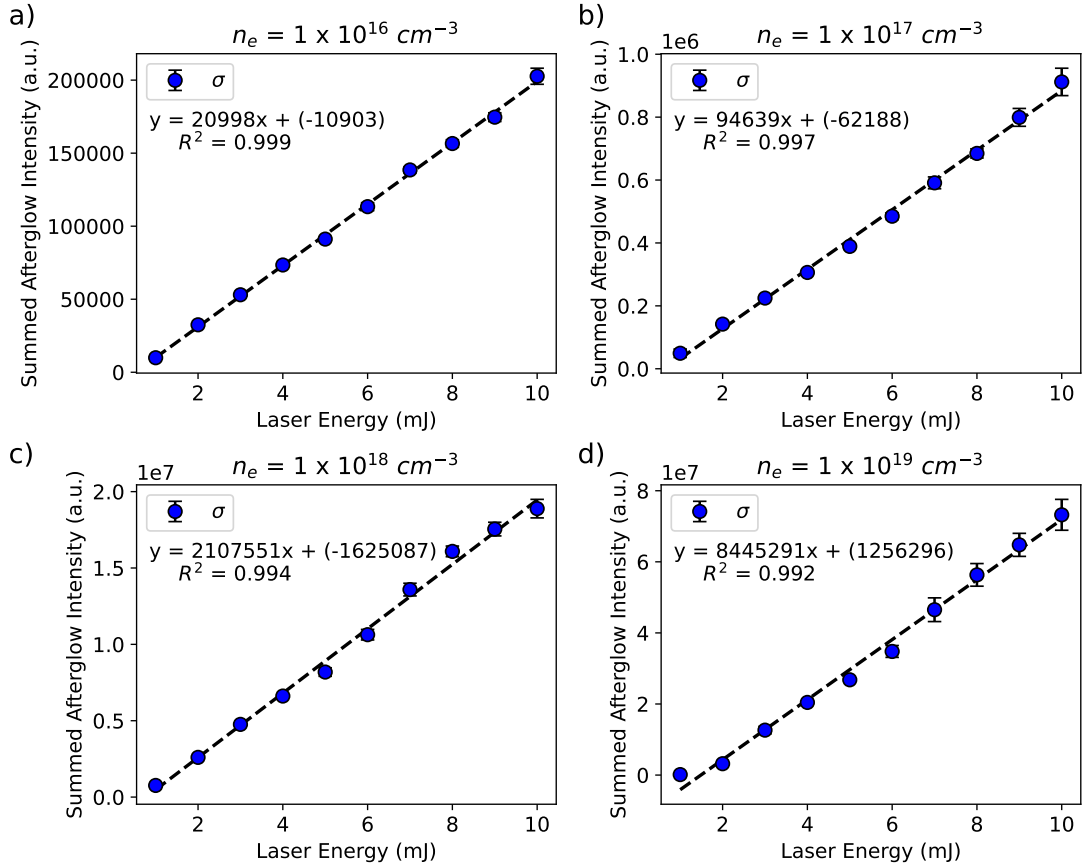


Figure 5.20: Wide range laser energy scan. Laser pulse's energy is varied and the corresponding summed afterglow signal's intensity is recorded (blue markers) at 4 different gas densities - $n =$ a) 10^{16} b) 10^{17} c) 10^{18} & d) 10^{19} cm^{-3} . Each marker represents the mean summed afterglow intensity for 50 data points. A linear regression model is applied to the data (dashed line) with the line of best fit equation and R^2 value displayed to quantify the proximity of our data to the fitted regression line. The error bars represent the standard deviation, σ .

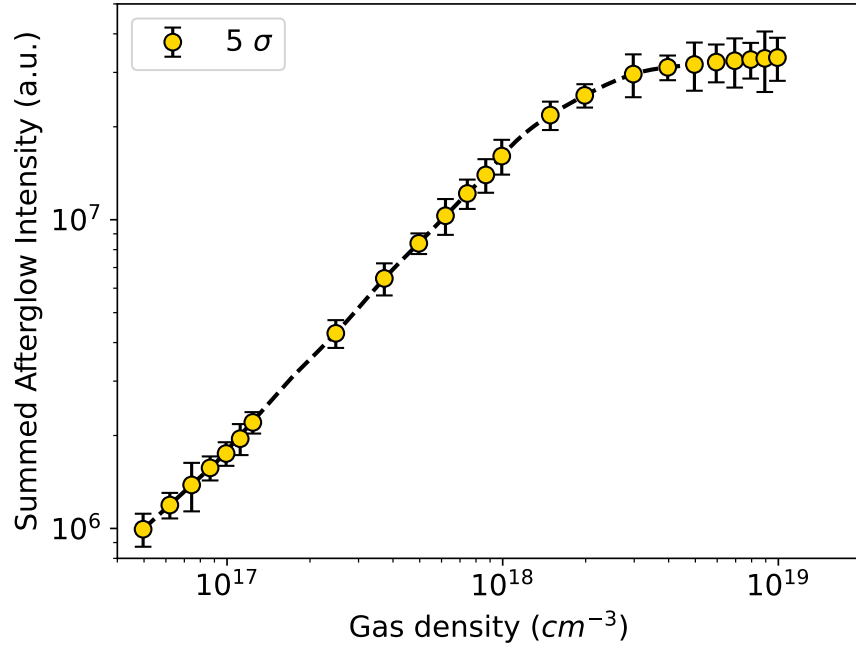


Figure 5.21: Wide range gas density scan. Gas density is varied between $5 \times 10^{16} \text{ cm}^{-3}$ and $1 \times 10^{19} \text{ cm}^{-3}$ in steps of $1 \times 10^{16} \text{ cm}^{-3}$ between $5 \times 10^{16} - 1 \times 10^{17} \text{ cm}^{-3}$, $1 \times 10^{17} \text{ cm}^{-3}$ between $1 \times 10^{17} - 1 \times 10^{18} \text{ cm}^{-3}$ and $1 \times 10^{18} \text{ cm}^{-3}$ between $1 \times 10^{18} - 1 \times 10^{19} \text{ cm}^{-3}$, with the corresponding summed afterglow signal's intensity recorded at each step by SideView CCD 2. Each marker represents the mean summed afterglow intensity for 50 data points. The error bars encompass 5 times the standard deviation, σ .

species is crucial as employing the diagnostic around this value will lead to non-linearities in the regression model. It is also important to note that the scaling identified here is representative of the afterglow emission visualised in one plane, the z-y plane in this case. If one was to image the entire volume of the plasma and subsequent afterglow light, the laser energy - summed afterglow intensity relationship is more likely to resemble a cubic function.

5.2.3 Gas density variation

Given its significance to the operation of an LPA, a wide ranging scan of the gas density was prioritised and performed in ‘flooded chamber’ mode, where the

gas density n was varied between $5 \times 10^{16} \text{ cm}^{-3}$ and $1 \times 10^{19} \text{ cm}^{-3}$ whilst keeping the laser parameters fixed at the nominal values - $\lambda_L = 800 \text{ nm}$, $W_L = 5.16 \text{ mJ}$ and $t_{FWHM} = 220 \text{ fs}$. The parameter range limits were selected due to the increasing uncertainty when measuring small changes in pressure at values below $5 \times 10^{16} \text{ cm}^{-3}$ and concerns about the amount of radiation being produced at values beyond $1 \times 10^{19} \text{ cm}^{-3}$, given the level of protection in the experimental lab. Figure 5.21 displays a log-log plot of the mean summed afterglow signal intensity for 50 data points, as recorded by SideView CCD 2 for each gas density value.

It is noticeable here that the afterglow signal is linear from roughly $5 \times 10^{16} \text{ cm}^{-3}$ to $1 \times 10^{18} \text{ cm}^{-3}$ but then starts to increasingly flatten and tail off as the gas density is further increased between $1 \times 10^{18} \text{ cm}^{-3}$ and $1 \times 10^{19} \text{ cm}^{-3}$, thus we can infer linearity between the gas density and afterglow signal intensity in certain ranges but not across the entire domain. Each afterglow camera was closely monitored between $1 \times 10^{18} \text{ cm}^{-3}$ and $1 \times 10^{19} \text{ cm}^{-3}$ to confirm the afterglow signal intensity roll off is a physical effect in the plasma, and not the result of CCD sensor saturation. The linear relationship displayed in the low gas density range is significant as it allows for the afterglow to be utilised as an in-situ method for measuring densities which are often inaccessible via interferometry. Furthermore, it is worth noting how minimal the standard deviation is for the afterglow signal intensity in these measurements, with the corresponding error bars having to cover 5σ instead of 1σ so that they are visible on the plots.

To gain greater insight into this parameter scan, the mean afterglow images recorded by SideView CCD 2 from 50 data points are displayed in Figure 5.22 for 6 gas density cases - 5×10^{16} , 1×10^{17} , 5×10^{17} , 1×10^{18} , 5×10^{18} and $1 \times 10^{19} \text{ cm}^{-3}$, with the laser's vacuum focal position along z highlighted via the white dashed line. From these images, we can see that the mean position of the afterglow signal along z , $\bar{A}$, calculated by averaging the location of the 50

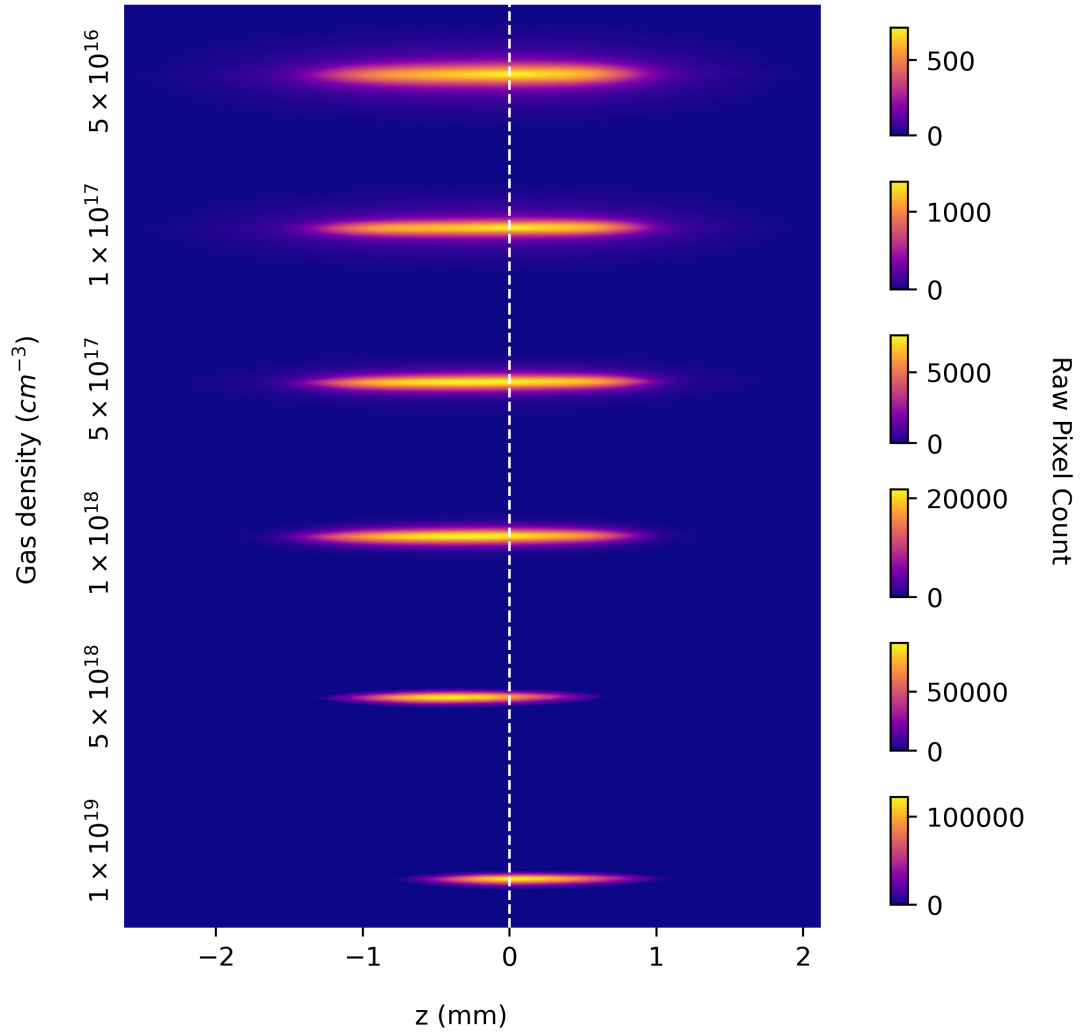


Figure 5.22: Afterglow Images - Gas density variation. Images of the mean afterglow signal recorded by SideView CCD 2 from 50 data points for 6 different gas density cases - 5×10^{16} , 1×10^{17} , 5×10^{17} , 1×10^{18} , 5×10^{18} & $1 \times 10^{19} \text{ cm}^{-3}$. The white dashed line represents the laser's vacuum focal position along z .

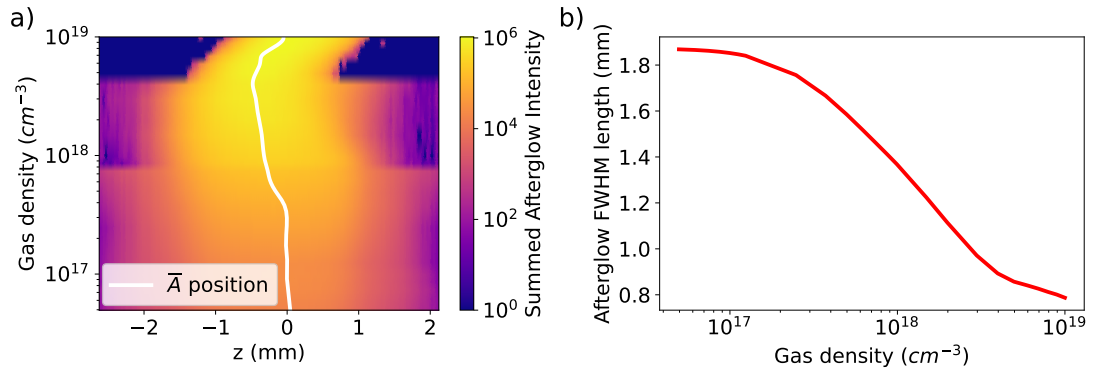


Figure 5.23: Wide range gas density scan afterglow trends. In subplot a), the mean afterglow image (as shown in Fig 5.10 a)) for each gas density has been summed along the y axis to create one afterglow signal intensity slice along z . These are then stacked on top of each other and plotted to visualise the afterglow signal's evolution along z as the gas density is increased. The white line signifies the mean position of the afterglow signal along z , $\bar{A}$. In subplot b), the red line represents the afterglow FWHM length as the gas density is increased.

highest intensity pixels on the camera chip, shifts substantially both towards and then away from the focusing lens. At the same time, the FWHM length of the afterglow signal decreases as the gas density is increased, despite the intensity of the laser pulse entering the vacuum chamber nominally remaining constant.

These two trends are laid out more clearly in Figure 5.23 a), where the mean afterglow image, recorded by SideView CCD 2 from 50 data points, for each gas density has been summed along the y axis and then stacked on top of each other to aid in visualising the signal's evolution along the z axis, and b). An OD = 0.2 ND filter was placed in front of the camera before recording the afterglow at densities beyond $4 \times 10^{18} \text{ cm}^{-3}$ due to the light becoming intense enough to cause camera pixels to saturate, which has resulted in the low intensity afterglow light on the far left and right hand side of the CCD chip being filtered out. Similarly, the exposure time of SideView CCD 2 was reduced between 7×10^{17} and $8 \times 10^{17} \text{ cm}^{-3}$ to prevent oversaturation, meaning some low intensity light at the outer regions of the focal plane was not detected in the neighbouring gas

density afterglow images.

5.3 Simulations

In an attempt to explain the features highlighted during the wide range gas density scan in Section 5.2.3, further FBPIC simulations were performed to build on the initial run in Section 5.1.4, in which the gas density was varied between $5 \times 10^{16} \text{ cm}^{-3}$ and $1 \times 10^{19} \text{ cm}^{-3}$ whilst keeping the laser parameters fixed at the nominal values - $\lambda_L = 800 \text{ nm}$, $W_L = 5.16 \text{ mJ}$ and $t_{FWHM} = 220 \text{ fs}$, thus mimicking the experiment carried out in Section 5.2.3. The setup parameters were kept the same as the initial run with the simulation box size = $345.5 \times 924.8 \text{ } \mu\text{m}$, $D_z = 5 \times 10^{-8} \text{ } \mu\text{m}$, $D_r = 4 \times 10^{-7} \text{ } \mu\text{m}$, $N_m = 2$, $p_{nz} = 1$, $p_{nr} = 2$ & $p_{nt} = 4$. As before, stationary helium atoms were added as the particle species, with each macroparticle capable of being ionised via the ADK ionisation model, and the laser pulse was set to propagate for 4 mm, with the focus set at the midway point, 2 mm.

What was immediately clear upon analysis is that due to the significant amount of energy contained within each pulse, we ionize well beyond the beam's Rayleigh length, as previously shown in Figure 5.9 a), which was calculated to be $z_R = 94 \text{ } \mu\text{m}$ using Equation 2.17, by $\approx 1.5 \text{ mm}$ each side of the focus for the linear, low gas density cases. This substantial amount of ionization gives rise to increased plasma absorption via inverse bremsstrahlung [260], where free electrons oscillating in the pulse's electric field absorb energy from the laser's photons, and pronounced ionization-induced defocusing (see Section 2.3.2) of the pulse as the gas density is increased, which causes the beam to diverge away from the propagation axis and its intensity to drop well below the diffraction limited value [22].

This is ultimately responsible for the afterglow's reduction in length and back-

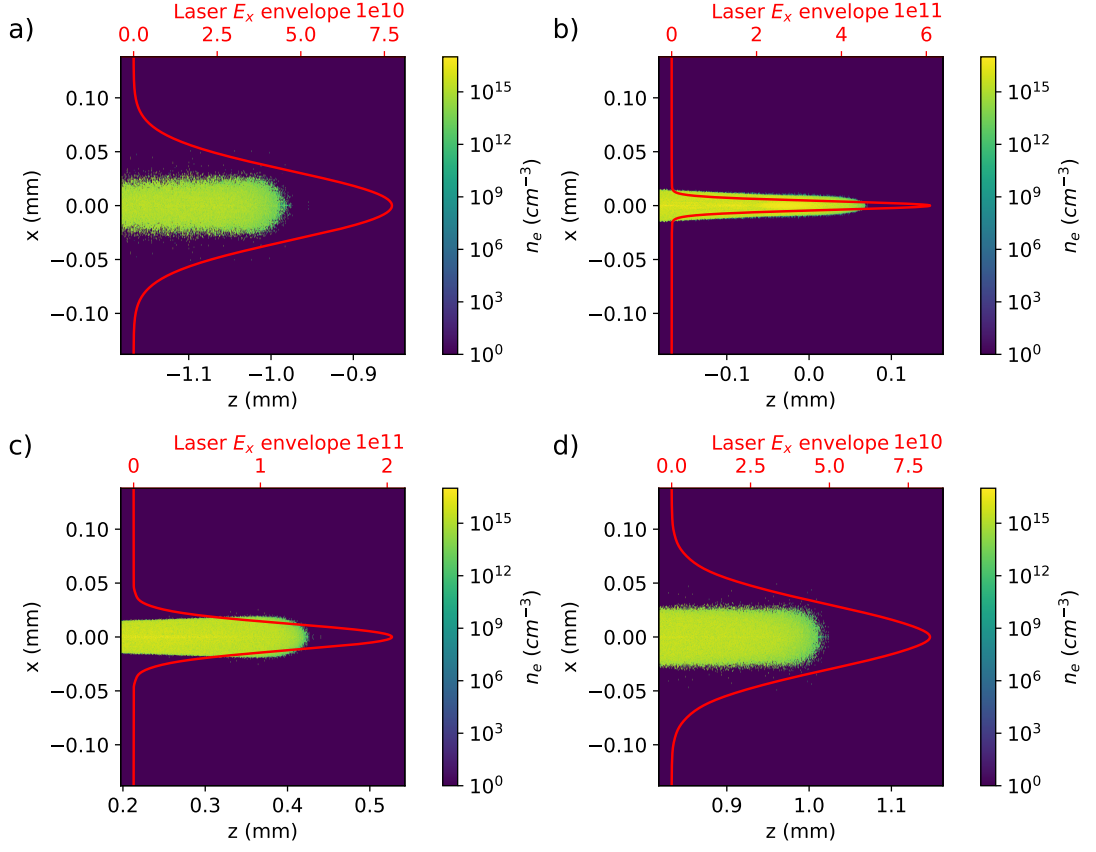


Figure 5.24: Linear laser-gas propagation and focusing case. FBPIC simulated laser pulse envelope and ionization progression along z for the linear, low gas density $5 \times 10^{16} \text{ cm}^{-3}$ case. The laser pulse's electric field envelope (red line) is overlaid on top of the laser ionized electron density and displayed at 4 different points along the propagation w.r.t. the focal position: a) -1.00 mm b) 0.00 mm c) 0.38 mm and d) 1.00 mm. Here the laser follows a traditional focus converging towards and then diverging away from the lens' focal point, with the electric field envelope and level of ionization mirrored on both sides of the focus.

ward movement along the propagation axis z towards the focusing lens shown in Figures 5.22 & 5.23, as increasing the gas density causes plasma absorption and ionization-induced defocusing to become more prominent, meaning a larger portion of the pulse's energy is lost to the plasma and the beam is deflected increasingly earlier along z . Therefore, the distance the laser is able to ionize post vacuum focal position becomes progressively smaller, leading to a reduction in the overall size and shift to the left in the mean position of both the plasma filament and subsequent afterglow signal. More precisely, the afterglow FWHM length drops from 1.85 mm to 1.03 mm and $\bar{A}$ moves to left by 0.42 mm between $n = 5 \times 10^{16} \text{ cm}^{-3}$ and $5 \times 10^{18} \text{ cm}^{-3}$.

A clear signature of ionization-induced defocusing taking place in the simulations was the laser starting to split into two distinct propagation modes [261] as it moved through the vacuum chamber - the portion of the beam traveling within the central plasma column that drops in intensity and the remainder of the pulse traveling outside it. This is demonstrated in Figure 5.25, which displays the laser pulse's electric field envelope (red line) and ionized electron density for the non-linear, high gas density $4 \times 10^{18} \text{ cm}^{-3}$ case, where the field envelope dips inward from its peak, increasingly deviating away from its Gaussian profile to form a ring-like structure [262]. As discussed in Section 2.3.2, this is caused by a positive radial index gradient created in the plasma by the intensity-dependence of the ionization, which results in the formation of a defocusing lens for the portion of the beam inside the central plasma column. For comparison, the laser pulse's electric field envelope (red line) and ionized electron density for the linear, low gas density $5 \times 10^{16} \text{ cm}^{-3}$ case is displayed in Fig 5.24, where the focused laser pulse propagates in a conventional manner, affirming that ionization-induced defocusing is contingent on the selected gas density.

For the high gas density cases where n exceeds $1 \times 10^{18} \text{ cm}^{-3}$, ionization-induced defocusing becomes severe enough to cause the laser intensity to drop

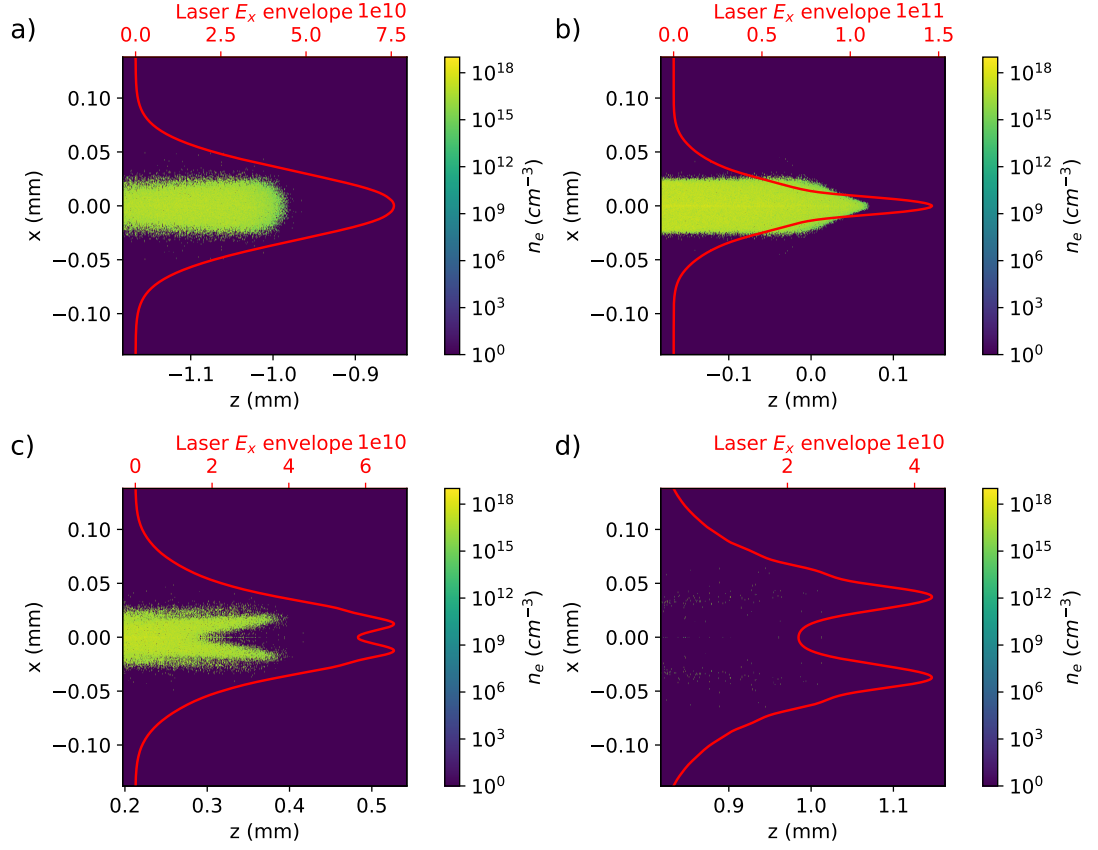


Figure 5.25: Ionization-induced defocusing laser-gas propagation and focusing case. FBPIC simulated laser pulse envelope and ionization progression along z for the non-linear, high gas density $4 \times 10^{18} \text{ cm}^{-3}$ case. The laser pulse's electric field envelope (red line) is overlaid on top of the laser ionized electron density and displayed at 4 different points along the propagation w.r.t. the focal position: a) -1.00 mm b) 0.00 mm c) 0.38 mm and d) 1.00 mm. In a) & b), the pulse profile is Gaussian and the laser steadily ionizes the background helium as it focuses down. In c), the Gaussian profile separates as a consequence of ionization-induced defocusing and the central portion of the laser's electric field envelope dips below the ionization threshold. In d), this defocusing effect has clearly become more pronounced with the entire pulse now incapable of ionizing the helium inside the vacuum chamber.

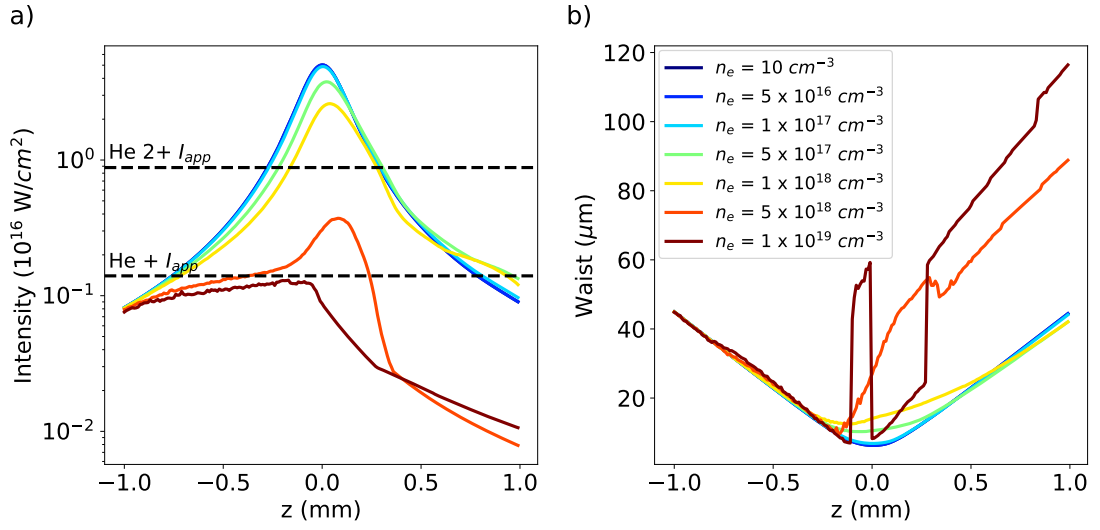


Figure 5.26: Simulated laser pulse intensity and waist evolution. The laser pulse a) intensity and b) waist evolution along the propagation, ± 1 mm from the focus, simulated for several gas density cases. A specific color was assigned to each gas density and displayed in the legend in b)). The dashed black lines in a) represent the He + & He 2+ appearance intensities, I_{app} , respectively.

below the Helium 2+ I_{app} threshold and so the pulse gradually becomes unable to ionize this second level ion species. This can be seen in Figure 5.26 where the FBPIC simulated laser pulse a) intensity and b) waist evolution along the propagation is displayed for $n = 10, 5 \times 10^{16}, 1 \times 10^{17}, 5 \times 10^{17}, 1 \times 10^{18}, 5 \times 10^{18}$ and $1 \times 10^{19} \text{ cm}^{-3}$ cases. These quantities are calculated by extracting both the intensity and spatial dimensions of the laser's field envelope at each simulation time step. The spike visible in b) in the $n = 1 \times 10^{19} \text{ cm}^{-3}$ case's laser waist (brown line) is a result of the field envelope becoming severely distorted (see Figure 5.28), causing large fluctuations and uncertainty in the value calculated by our analysis function, that computes the RMS radius weighted by the transverse profile.

This drop below the Helium 2+ I_{app} threshold coincides with and is likely responsible for the 'levelling off' we see at the same gas density values in Fig 5.21, though this assumption is difficult to unpack as there are competing effects

taking place here. Generally, higher gas density means there are more atoms present that can be ionized and so the level of emission would be expected to increase. However as previously discussed, higher gas density also promotes more extreme ionization-induced defocusing causing the laser intensity to drop and both shorten the afterglow signal in z and prevent He 2+ ionization in this case, which will result in less light being emitted from the plasma.

Advantageously, the afterglow was discovered to be reactive to these ionization-induced defocusing induced changes in the laser pulse's intensity and waist as shown in Figure 5.27. More specifically, we see in subplot a) that the shift in the afterglow's mean position along z can be used to determine the laser's focus/minimum waist position along z , which as previously discussed moves towards the left as the density is increased due to the pulse being deflected away from the propagation axis more quickly. Although the directional trend is consistent across both quantities, the magnitude of the mean afterglow positional shift was found to be approximately three times greater than that of the minimum waist position as it also encompasses the significant plasma absorption via inverse bremsstrahlung taking place here. Oscillations in the linear trend line displayed in subplot b) can be derived from the natural discrepancy between the FBPIC simulated parameters and those present on-target in the experiment, which ultimately leads to some of the subtle positional changes occurring at slightly different gas density values.

Furthermore in subplot c), the FWHM length of the afterglow is demonstrated to decrease in accordance with the laser pulse's peak intensity I_0 for our experimental case. Although it was shown in Figure 5.19 that the afterglow FWHM length does not vary linearly with the laser energy/intensity but with the He+ I_{app} position along z , here the laser's energy as it enters the vacuum chamber stays unchanged and instead the afterglow length is able to reflect the non-linear reduction in the pulse's peak intensity, as seen in Figure 5.26, within the gas/plasma as a

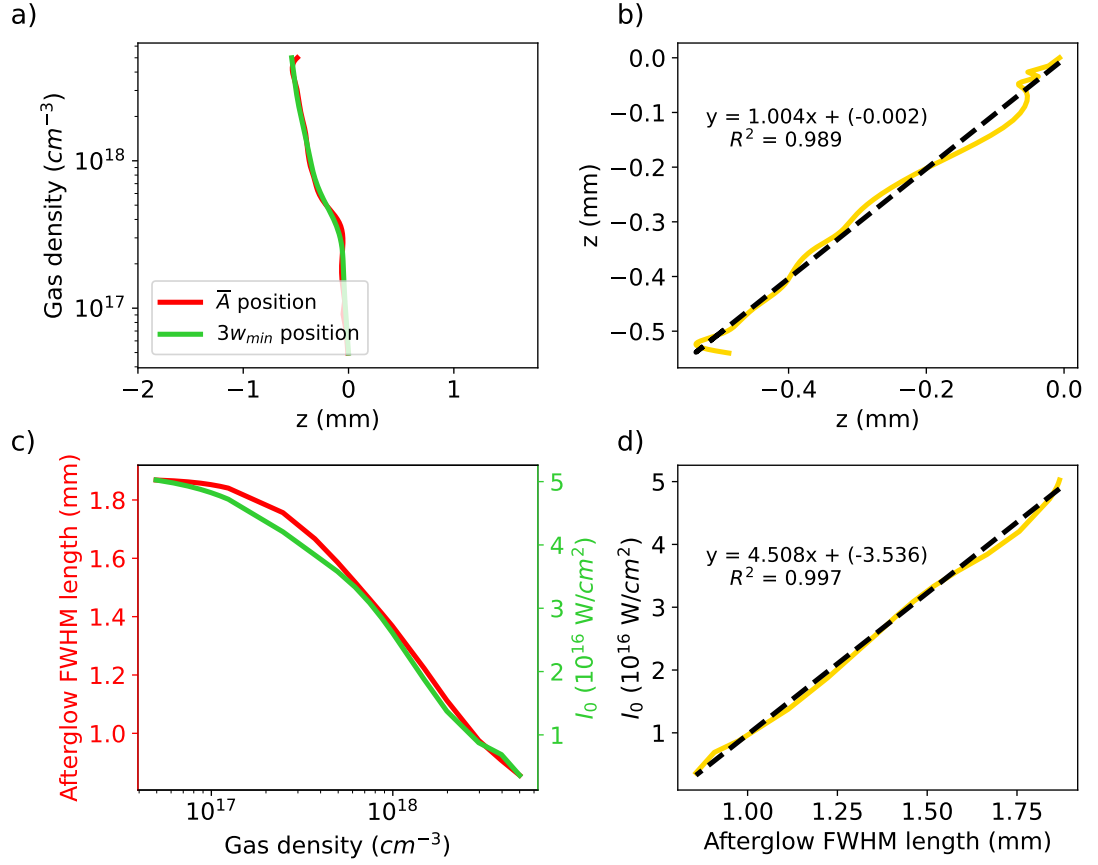


Figure 5.27: Ionization induced defocusing afterglow metrology. Afterglow signal is found to be receptive to changes in the laser's minimum waist, w_{min} , position along z and peak intensity, I_0 . Subplot a) displays the mean afterglow position, $\bar{A}$, and the w_{min} position multiplied by a factor of 3 during the gas density scan, with b) presenting these two quantities plotted against each other with a linear regression model applied to the data. Similarly, c) shows the afterglow FWHM length and I_0 during the gas density scan, with d) presenting these two quantities plotted against each other with a linear regression model applied to the data.

result of ionization-induced defocusing. The afterglow can therefore be deployed to diagnose the actual intensity we are receiving on-target, accounting for the presence of ionization-induced defocusing, which is a major first step in terms of optimizing of a laser-plasma accelerator to consistently deliver the desired output electron beam parameters.

Notably, the very high density cases (between $5 \times 10^{18} \text{ cm}^{-3}$ & $1 \times 10^{19} \text{ cm}^{-3}$) where the mean position of the afterglow signal along z can be visualised moving back towards the vacuum focus in Figure 5.23, have been omitted from the trend lines in Figure 5.27. At these values, it became apparent that self-focusing had begun to become a feature due to the gas density dependence on the critical power for self-focusing of a laser pulse in a gas $P_{crit} = 3.77\lambda_L^2/8\pi n_0 n_2$ [263] where λ_L is the laser wavelength and n_0 and n_2 characterize the intensity-dependent refractive index $n = n_0 + n_2 I$. More specifically, moving into this density range caused the critical power value to drop by an order of magnitude by way of the exponential increase in the non-linear refractive index, n_2 [264].

This triggers the scenario displayed in Figure 5.28, in which the pulse focuses down and undergoes ionization-induced defocusing to create a spatial ring profile as before, however after further propagation, the trailing portion of the central pulse reforms spatially due to self-focusing in a process termed ‘dynamic spatial replenishment’ by Mlejnek *et al.* [265]. This means that a second focus is generated beyond the conventional vacuum focal position, allowing for ionization to occur further along z than in the previous gas density cases and subsequently shifting the mean afterglow signal position along z back towards the right.

As is apparent from Fig 5.28 , the pulse is highly distorted by the laser-gas interaction at these density values and the empirical model, based on experimental and simulated data, laid out in Fig 5.27 for diagnosing ionization-induced defocusing breaks down as the afterglow can no longer be utilised to monitor the minimum laser waist position in the presence of multiple foci. Regardless, this

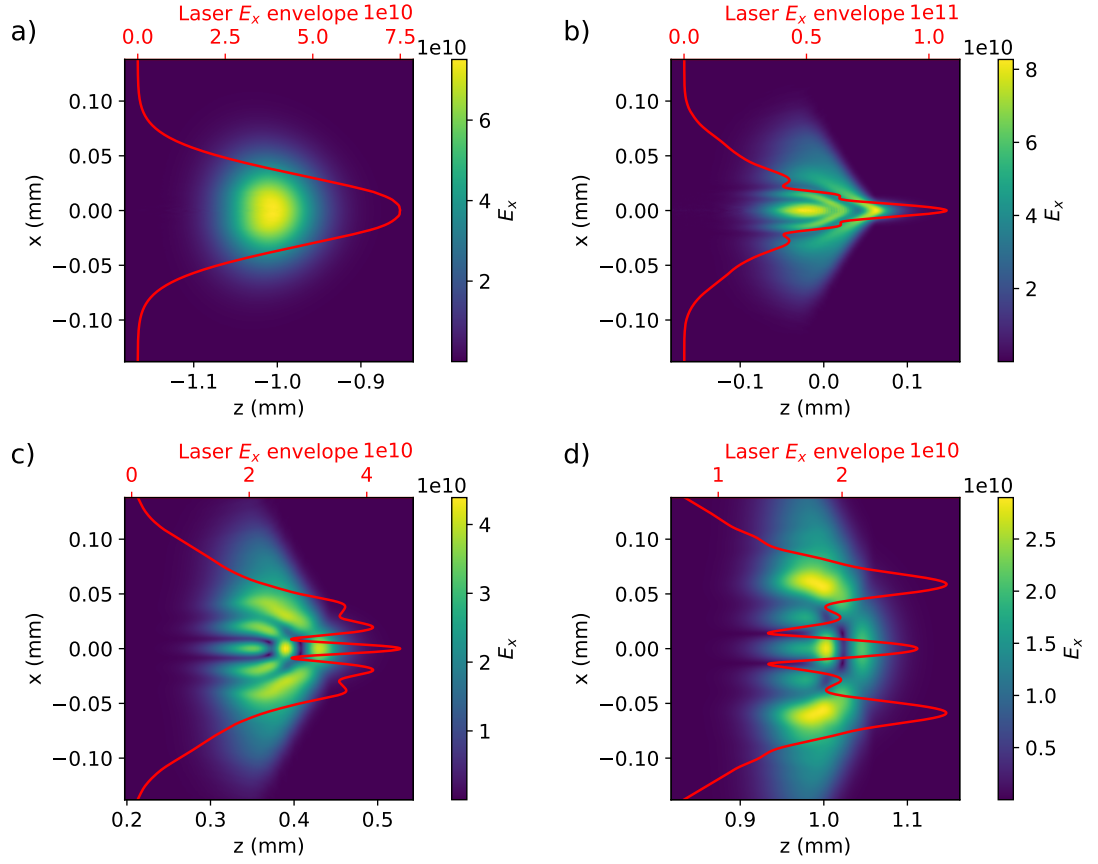


Figure 5.28: Dynamic spatial replenishment propagation case. The laser pulse envelope and E_x field progression along z for the highly non-linear, high gas density $1 \times 10^{19} \text{ cm}^{-3}$ case. The laser pulse's electric field envelope (red line) is overlaid on top of the laser's E_x field, the magnitude of the electric force vector acting along the x -axis, and displayed at 4 different points along the propagation w.r.t. the focal position: a) -1.00 mm b) 0.00 mm c) 0.38 mm and d) 1.00 mm. In subplot b), the leading portion of the pulse can be seen focusing down acutely. Subsequently, the laser defocuses into the signature ionization-induced defocusing ring-like structure before, as a result of self-focusing, the trailing central portion of the pulse reforms spatially as displayed in c) in a process termed 'dynamic spatial replenishment'. This allows for ionization to occur past the vacuum focus in z before the pulse starts to decay due to non-linear absorption as shown in d).

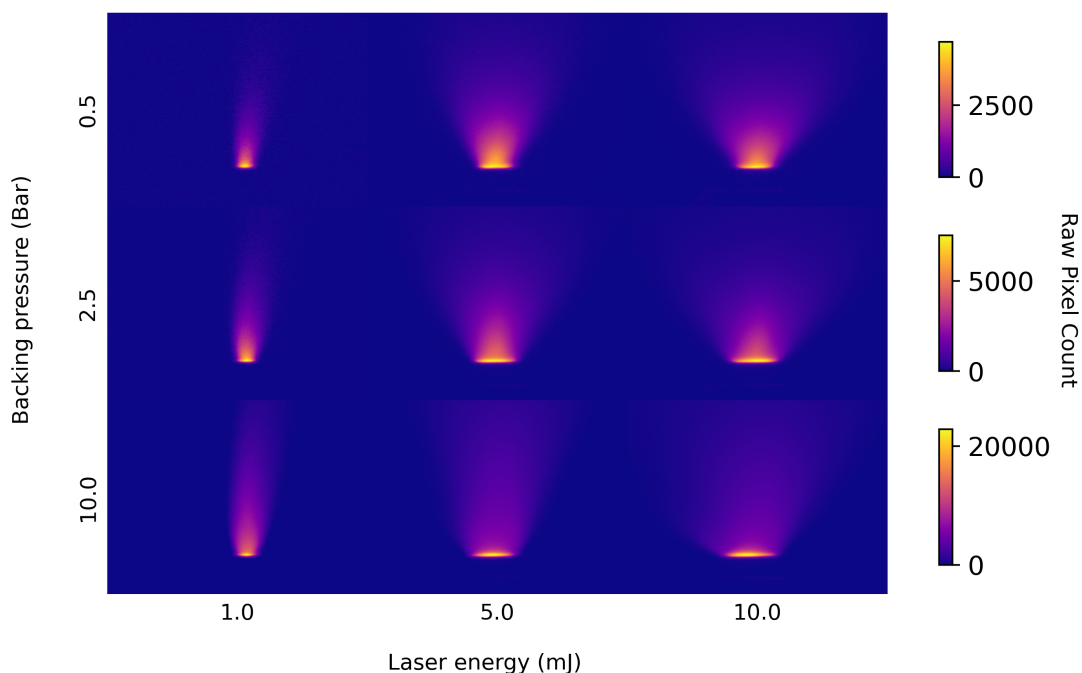


Figure 5.29: Afterglow Images - Gas jet mode. Images of the mean afterglow signal in gas jet mode recorded by SideView CCD 2 from 50 data points for 3 different laser energies - 1, 5 and 10 mJ, and 3 different backing pressures - 0.5, 2.5 and 10.0 Bar.

model is still robust and relevant over a gas density range spanning 3 orders of magnitude, within which the majority of LPA regimes lie.

5.4 Experimental results - Gas jet mode

After establishing the afterglow's relationship with the gas density in flooded chamber mode, there was a desire to switch to gas jet mode, given its status as the increasingly preferred method by which gas is supplied in an LPA experiment. Unlike in flooded chamber mode where the laser interacts with a constant gas reservoir, here gas is injected from a nozzle creating a highly localised, well-defined gas cloud in a vacuum. As discussed in Section 5.1.3, a 3 mm SourceLAB SL-NOZ I nozzle was attached to the valve, which has been successfully employed

in hybrid LPWFA experiments [77, 266]. Mean images of the afterglow signal in gas jet mode recorded by SideView CCD 2 from 50 data points are displayed in Figure 5.29 for 3 different laser energies - 1, 5 and 10 mJ, and 3 different backing pressures - 0.5, 2.5 and 10.0 Bar. In these images, we see a pronounced ‘spray’ feature where the afterglow signal appears to resemble the gas flow from the supersonic nozzle. Here, the helium ions move upwards after being ionised and the recombination/relaxation events and associated light is spread in the direction of the gas flow. The addition of a gated camera in place of SideView CCD 2 would allow this upward ion flow to be observed in real time so that its velocity could be determined as well as its correlation with the gas flow.

The primary aim in this experimental mode was to explore the afterglow signal’s ability to reconstruct the gas density profile of a particular nozzle. To do this, the gas jet was mounted on a 3-axis motorized mechanical stage and moved transversely with respect to the laser’s focus (along the x axis) in small increments whilst capturing the corresponding afterglow at each step using TopView CCD 3. Movement in this direction ensures the laser’s intensity profile remains constant across the width of the jet, so the only quantity varying is the gas profile. This scan was performed at our nominal laser parameters - $\lambda_L = 800$ nm, $W_L = 5.16$ mJ, $t_{FWHM} = 220$ fs and $I_0 \approx 6 \times 10^{16}$ W/cm², at 4 different backing pressures - 0.5, 2.5, 4.0 & 10.0 Bar, with the gas jet moved in increments of 0.2 mm, ± 2 mm either side of the nozzle centre.

Mean images of the afterglow signal during this gas jet transverse movement scan recorded by TopView CCD 3 from 50 data points are displayed in Figure 5.30 for the 2.5 Bar case at 5 different transverse positions - -0.9, -0.5, -0.1, 0.3 and 0.7 mm. Here we see the afterglow reflect that the focused laser pulse is sampling different sections of the gas density profile, with the light emission decreasing from the centre towards the nozzle edges. In order to properly benchmark these results, simulations of the gas flow through the nozzle were performed using Ansys

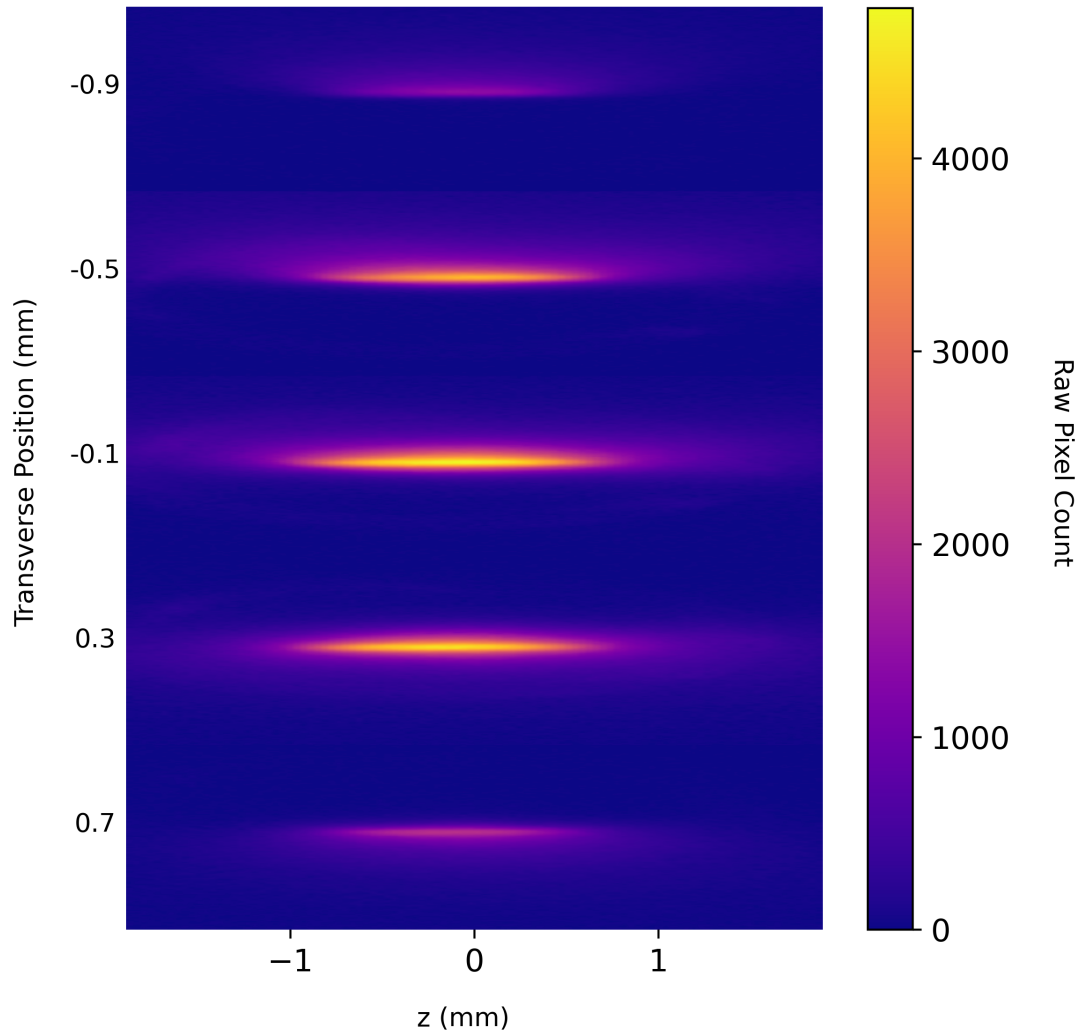


Figure 5.30: Afterglow Images - Gas jet transverse movement variation. Images of the mean afterglow signal during the gas jet transverse movement scan recorded by TopView CCD 3 from 50 data points for the 2.5 Bar case at 5 different transverse positions - -0.9, -0.5, -0.1, 0.3 and 0.7 mm.

Fluent to deliver density contours of the gas. A line-out of these contours was then taken and plotted alongside the summed afterglow intensity measurements for each backing pressure case, as seen in Fig 5.31.

On display here is the afterglow's ability to effectively re-create the nozzle's gas density profile with a high degree of spatial sensitivity. For the 4.0 & 10.0 Bar cases shown in Fig 5.31 c) & d), the afterglow light does struggle to resolve the simulated gas profile's u-shaped top, however this is believed to be a simulation artefact rather than a genuine physical feature. It is important to remember that the actual gas profile could be subtly different to the Ansys Fluent profile as the simulation does not consider the geometry of the valve utilised, the piping and the jet's opening time. In future measurements, this could be directly tested by moving the gas jet in smaller increments around the wings of the gas profile to better resolve the large-scale density variations that occur over very fine distances in these regions.

The same non-linearities that occur at high gas densities in a 'flooded chamber' will also be present in this operational mode and affect the afterglow's ability to reconstruct the gas profile. Knowing this and studying the 0.5 & 2.5 Bar cases, we can assert that the technique performs better at lower densities which is advantageously the inverse of conventional measurement methods such as interferometry that are known to become more inaccurate as the gas density is decreased. This has particular significance for correctly diagnosing the low density wings of the gas jet which can directly impact the laser-plasma-acceleration process via plasma lensing of the laser pulse on the upramp, and emittance preservation of generated electrons beams on the downramp.

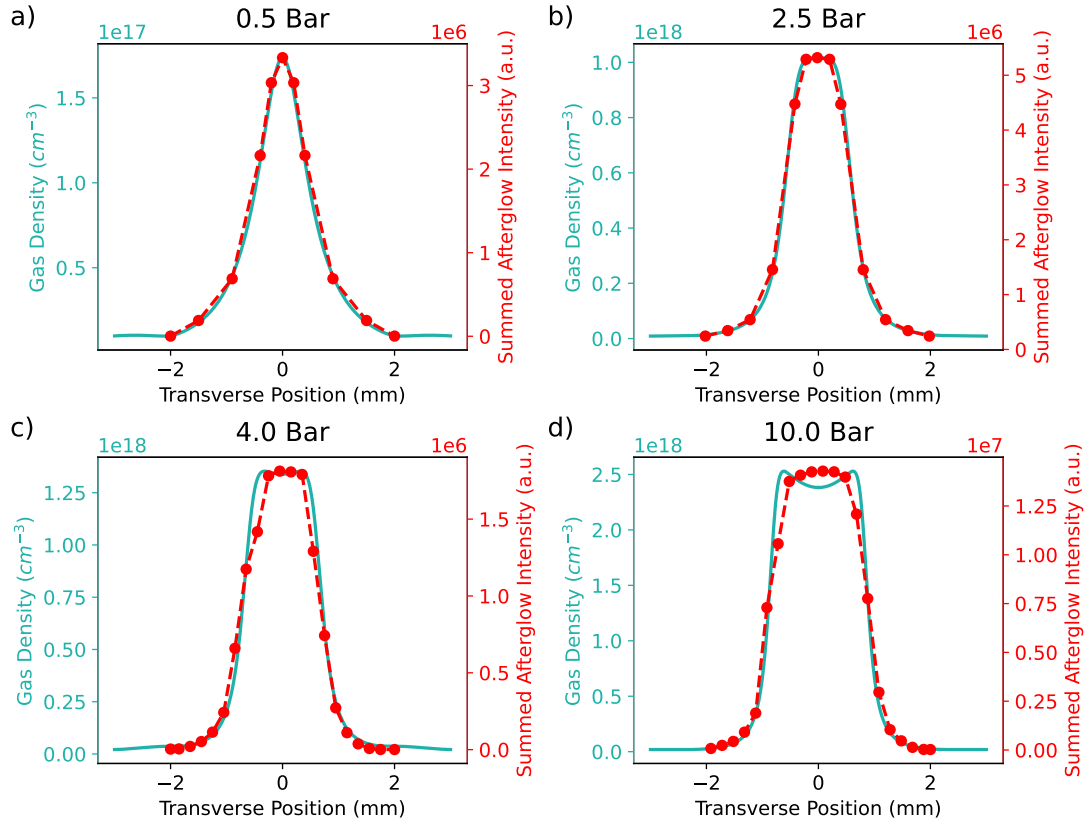


Figure 5.31: Gas nozzle profile reconstruction. Comparison between the Ansys Fluent simulated 1D gas profile and the experimentally captured afterglow 1D gas profile, reconstructed by incrementally moving the gas jet transversally w.r.t. the laser's focus and recording the afterglow signal at each step, of a 3 mm SourceLAB SL-NOZ I nozzle at 4 different backing pressures: a) 0.5 Bar b) 2.5 Bar c) 4.0 Bar d) 10.0 Bar. The teal colored line corresponds to the Ansys fluent generated profile, simulated ± 3 mm each side of the nozzle centre, and the red line represents the experimental afterglow signal profile, catalogued in 0.2 mm increments, ± 2 mm either side of the nozzle centre.

5.5 Summary

The work summarised in this chapter represents a significant step forward in terms of establishing the capabilities of the plasma afterglow signal as a diagnostic for laser-plasma accelerators. This was aided by creating a well controlled experimental setup with an extremely stable plasma seed as oppose to previous studies where the afterglow has been recorded as a parasitic measurement with less stable lasers.

The correlation between the afterglow and several key experimental parameters was presented, alongside a theoretical justification for each trend and discussions centred on important considerations when employing the technique to measure a specific quantity. This coupled with the inherent advantages of the measurement, its non-destructive nature and the simplicity of the required setup substantiates the afterglow as an increasingly attractive approach for diagnosing an accelerator in real time, especially when experimentally constrained.

Furthermore, the existing scope of the plasma afterglow was expanded by demonstrating its capacity to emulate a non-linear propagation effect in the form of ionization-induced defocusing experienced by the laser as it traverses through the gas. Being able to resolve these complex laser-plasma dynamics is unique and sets the technique apart from comparable diagnostics.

Looking more broadly within the field of research, this technique lends itself as both an intensity and spatial diagnostic for laser wakefield accelerator injection schemes like self-truncated ionization injection [71, 183], where precise control over the laser's peak intensity and focus evolution is required to achieve the desired electron beam spectra. Similarly, injection schemes in plasma accelerators that utilize accessory laser pulses [93, 191], like the method discussed in Chapter 4, are also prime candidates to employ afterglow to aid in clarifying the status of a focusing laser beam whilst it interacts with complex plasma based structures

like a wakefield.

It is important to acknowledge that in a real-life scenario, multiple experimental parameters are likely to jitter simultaneously and it will be difficult to decouple these by solely studying the afterglow signal. Engineering optimization of the setup in question is therefore still required in an attempt to minimize any significant jitters. However, even if numerous individual fluctuations are present, the afterglow is still an effective compound measurement and can be used for instance to return the accelerator to its designated optimized state.

For future studies, it would be interesting to operate in a laser-plasma interaction regime where self-focusing is the dominant non-linear process to distinguish whether the afterglow can also be used to represent this effect. This would be particularly relevant as many laser-plasma accelerator setups are now actively utilising self-focusing to their advantage to, for instance, ensure the desired laser intensity is realised at roughly the centre of the gas target [184].

Another experimental objective should be the formulation of a more complete model linking the initial energy transfer from the ionizing laser pulse to the eventual relaxation and recombination emission light, with each intermediate step more adequately explained and understood. This is a difficult task due to the variance in time scales over which these processes occur and the resolution required to resolved the physics in question, however work is already underway in this area by [48] in an attempt to realise this.

Finally, utilising a kHz repetition rate laser as the afterglow source leads to the accrual of extremely large datasets, opening up the possibility of applying machine learning to the afterglow signal, a concept which has already been deployed in laser plasma accelerators [267] to improve the beam quality and stability of an accelerated electron bunch [268]. As in [268], Bayesian Optimization seems like the ideal candidate to convert the afterglow into an online/live diagnostic with an active feedback loop, able to extract multiple laser & gas parameters

from the signal and then adjust these quantities accordingly to optimize the accelerator. Equally, a neural network, specifically a variational autoencoder based early warning system where the model receives images of the afterglow light as an input, would be an appealing technique to explore [269].

Chapter 6

Summary and Outlook

In summary, three key aspects of plasma-based particles accelerators have been investigated over the course of this project with the aim of progressing plasma accelerators towards fully deployable, application-grade machines. In Chapter 3, experimental testing and verification of a laser-plasma accelerator module optimised for kHz laser operation is presented, alongside demonstrating that kHz laser-plasma accelerators can create electron fluxes with properties analogous to those encountered in space radiation environments. In Chapter 4, we outline how laser-plasma accelerator output beam parameters can be effectively stabilised against several experimental fluctuations in a secondary beam-driven plasma photocathode-based wakefield stage. In Chapter 5, we establish a live non-intrusive accelerator diagnostic, plasma afterglow, by correlating its relationship with key accelerator parameters, thus enabling it to be leveraged as a method for measuring historically difficult to diagnose quantities and effects at the interaction point. Outlooks pertinent to each individual topic are presented in the summary of their corresponding chapters in Sections 3.1.3, 3.2.7, 4.8 and 5.5.

In closing, we reflect on the areas of convergence between the three primary themes of this study and hypothesize on prospective pathways for their increased

integration. As discussed in Chapter 3, high repetition rate lasers are presently limited to comparatively low pulse energies due to concerns surrounding thermal load, gain-medium damage thresholds and non-linear effects in the amplifier chain [270, 271]. However, innovative methods such as Optical Parametric Chirped Pulse Amplification (OPCPA) [272] and Coherent Beam Combining (CBC) [273] may allow for kHz lasers to achieve energies ≥ 200 mJ in the next decade [274]. This would open up the possibility of using output beams from kHz LWFAs to drive hybrid LPWFA setups, akin to the one outlined in Chapter 4. An all optical kilohertz rate plasma photocathode wakefield accelerator has the potential to be transformative in terms of its ability to produce electrons orders of magnitude brighter, stronger and higher frequency than state of the art.

It is generally acknowledged that next-generation plasma accelerators, particularly those targeting high-energy physics, will require operation in a multi-stage architecture [237]. Realising these staged accelerators constitutes one of the foremost technical challenges in the field of plasma accelerator research. The configuration described in Chapter 4 is well suited for deployment as the first stage of a multi-stage plasma accelerator, owing to its capacity to deliver electron bunches of superior quality relative to conventional sources within a markedly reduced spatial envelope. Moreover, the plasma afterglow diagnostic discussed in Chapter 5 is an attractive option for straightforwardly monitoring the stability of various key parameters at each accelerator stage. Crucially, the technique has now been shown to be capable of measuring the density profile of a plasma source (see Section 5.4), as well as the alignment of both laser pulses and electron beams with the respect to it [47, 275]. These characteristics will take on heightened importance in a staged plasma accelerator, where emittance preservation of the accelerated beam between each step is vital.

Adopting a more holistic viewpoint, future efforts should be directed towards increasing the electro-optical efficiency of laser driven plasma accelerators. Most

modern LPA's are powered by Ti:Sa lasers which are extremely inefficient, $\approx 2.5\%$, as their operation is typically dependant on multiple pump lasers, each with their own significant individual losses [276]. However, laser systems based on emerging technologies such as diode-pumped fiber and thin disk Yb:YAG lasers [277, 278] are considerably more efficient, 10-20% & 25-35% respectively, and capable of operating at hundreds of kHz. With such lasers likely to become commercially available in the next few years, employing them to drive a plasma accelerator would significantly increase the usability and wall-plug efficiency of the device.

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