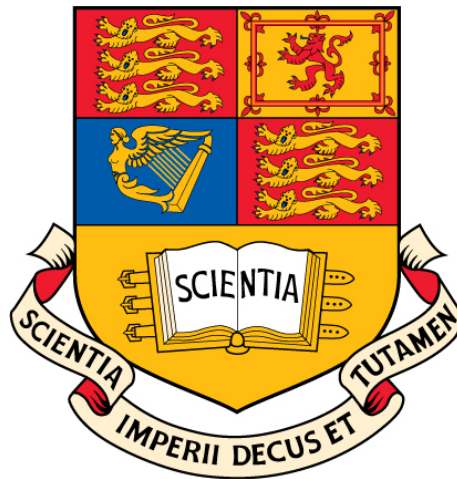


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Creating and Probing Warm Dense Matter and High Energy Density Blast Waves

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“ Every animal can be drawn as a series of monkey nuts ”

Muriel Kenny (1918-2009)

Abstract

Non-ideal plasmas are difficult to model numerically and achieve experimentally. Here I focus on reaching the warm dense matter (WDM) regime at one end of the energy scale and the high energy density (HED) regime at the other. Both regimes access common, yet hard to observe, features in planetary core physics and astrophysics respectively.

We have implemented several experiments developing both ionic, isochoric heating of matter to HED states, and probing the release isentrope through the WDM regime. Using sub-picosecond, ~ 100 J laser pulses, ion beams were produced in a thin foil interaction and then used to rapidly and uniformly heat ~ 20 μm wires to ~ 15 eV temperatures before hydrodynamic expansion could occur. Thus the partially ionised, strongly coupled WDM regime was accessed and then explored using streaked XUV and UV imaging of the expansion into vacuum, together with measurement of the ion spectra deposited in the sample. The calibrated ion spectra indicates the energy deposition occurs over ~ 400 ps and as such further investigation of the stopping power of warm dense matter is required.

The high absorption of these laser pulses in cluster gases ($> 90\%$) was used to create HED blast waves. Under hydrodynamic scaling laws these blast waves can be used to model processes occurring within astrophysical phenomena. Schlieren, interferometry and ion probing methods have been used to spatially and temporally resolve the blast wave profiles. We describe high energy scaling experiments in the search for the thermal cooling instability, in which radiation from the blast wave modifies both its density profile and propagation dynamics. A unique perpendicular, time delayed colliding blast wave geometry has been used to identify up-stream characteristics of the front propagation and to pre-heat material leading to shock front acceleration. We show that ions of an energy greater than ~ 200 keV are required to probe the transient electromagnetic fields believed to occur at the shock front.

The development and characterisation of low energy (0.1 - 10 MeV) proton beams was necessary for studies of both non-ideal plasma regimes, particularly for probing high density media and transient electromagnetic fields. The application of ions to recover stopping power is discussed for WDM and HED plasmas.

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The administration and technical help that supported my PhD cannot be underrated. Judith and Sara how you put up with us I’ll never know - although I have a suspicion

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If you have read this far you should probably carry on.

Fundamental Physical Constants

Physical Quantity	Symbol	Value (SI)
Boltzmann constant	k_B	$1.38 \times 10^{-23} \text{ J K}^{-1}$
Elementary charge	e	$1.60 \times 10^{-19} \text{ C}$
Electron mass	m_e	$9.11 \times 10^{-31} \text{ kg}$
Proton mass	m_p	$1.67 \times 10^{-27} \text{ kg}$
Speed of light in vacuum	c	$3.00 \times 10^8 \text{ m s}^{-1}$
Permittivity of free space	ϵ_0	$8.85 \times 10^{-12} \text{ F m}^{-1}$
Permeability of free space	μ_0	$4\pi \times 10^{-7} \text{ N A}^{-2}$
Stefan-Boltzmann constant	σ_{SB}	$5.67 \times 10^{-8} \text{ W m}^{-2} \text{ K}^{-4}$
Planck constant	h	$6.63 \times 10^{-34} \text{ Js}$

Contents

Abstract	3
Acknowledgements	4
Fundamental Physical Constants	6
List of Figures	11
List of Tables	15
1 Introduction	16
1.1 Contributors and Publications	17
2 Theoretical Background	19
2.1 Beyond Condensed Matter	20
2.1.1 Definition of Warm Dense Matter	20
2.1.2 The Equation of State for Warm Dense Matter	21
2.1.3 Methods for obtaining warm dense matter in the laboratory	22
2.1.4 Definition of a plasma	23
2.1.5 Modelling plasma motion	25
2.2 Laser-Matter Interactions	27
2.2.1 Ionisation Mechanisms	27
2.2.2 The Ponderomotive Force on an Electron in a Laser Field	28
2.2.3 Different Target Media for Laser Interaction Experiments	30
2.2.4 Energy Transfer Mechanisms in an Overdense Plasma	32
2.2.5 Laser Driven Ion Acceleration from Solid Targets	34

2.2.6	Electron and Ion Transport and Stopping in Matter	37
2.2.7	Laser-Cluster Interaction at High Intensities	38
2.3	Shocks in Plasmas	40
2.3.1	Shock Waves and Blast Waves	40
2.3.2	Astrophysical Shocks	41
2.3.3	Requirements for Laboratory Astrophysics Scaling	42
2.3.4	Producing an Optically Thin Radiative Blast wave in a Plasma . . .	44
2.4	Summary	46
3	Instrumentation and Experimental Methods	47
3.1	High Power Lasers	48
3.1.1	Chirped pulse amplification	51
3.1.2	Imperial College Nd:Glass laser	53
3.1.3	Optical Parametric Chirped Pulse Amplification	56
3.1.4	The Vulcan laser in target area west at the Rutherford Appleton Laboratory	57
3.2	Ion Diagnostics	59
3.2.1	Materials that detect ions	59
3.2.2	Detector stacks	64
3.2.3	Low energy magnetic spectrometer calibration	65
3.2.4	Thomson spectrometer	68
3.2.5	Ionography	68
3.3	Optical Diagnostics	69
3.3.1	Schlieren Imaging	69
3.3.2	Streaked Schlieren Imaging	70
3.3.3	Interferometry	71
3.4	Cluster Production	75
3.5	Summary	77
4	Proton Heating Towards Warm Dense Matter	78
4.1	Measurements for Equation of State Determination	79
4.2	Experimental Setup for Heating and Probing	79

4.3	Results	81
4.3.1	Time Resolved Density	82
4.3.2	Deposited Energy Determination from Ion Spectra	86
4.3.3	Uniformity of Heating from Streaked UV Emission	90
4.4	Discussion	93
4.5	Summary	99
5	High Energy Density Blast Waves in Clusters	100
5.1	Probing Upstream Material via Two Beam Geometries	101
5.1.1	Previous Work	101
5.1.2	Experimental Setup	102
5.1.3	Results	107
5.1.4	Discussion	109
5.1.5	Summary	113
5.2	Proton Probing of a Cluster Gas Blast Wave	114
5.3	High Energy Scaling of Single Blast Waves	117
5.3.1	Experimental Setup	118
5.3.2	Absorption Results	118
5.3.3	Supercritical Shocks	122
5.3.4	Fine Structure in Shocks	123
5.4	Summary	126
6	Proton and Stopping Power Characterisation	127
6.1	Full Proton Beam Characterisation	128
6.1.1	Low Energy Proton Spectra at Different Laser Energies	128
6.1.2	MAGPIE Proton Self-Emission from a Radial Wire Array	130
6.1.3	Virtual Proton Source Driven by a Low Energy Laser	131
6.1.4	Energy Dependent Divergence	133
6.1.5	Low Energy Transverse Emittance	135
6.2	Determining the Stopping Power	137
6.2.1	Methods of Characterisation Using Laser Generated Protons	138
6.2.2	Methods Utilising Mono-Energetic Proton Beams	141

6.3	Summary	143
7	Conclusions and Future Work	145
7.1	Advanced Optical Probing	147
7.2	Annular Shocks for Jet Production	150
7.3	New Laser Capability: Cerberus Laser	152
7.4	Concluding Remarks	153
A	Derivation of Magnetic Spectrometer Error	155
	References	158

List of Figures

2.1	Aluminium temperature-density phase diagram showing coupling and degeneracy	20
2.2	Aluminium temperature-density phase diagram showing equation of state models	22
2.3	Diagram of optical ionisation processes	28
2.4	Diagram of different target media	31
2.5	Diagram of ion acceleration mechanism	36
2.6	Radiochromic film exposed to protons	37
2.7	Nanoplasma model output showing a transient resonance	39
2.8	Diagram showing Drake's classification of astrophysical plasma opacity . . .	42
2.9	Graph showing the opacity region ionised Xenon gas occupies	45
3.1	Diagram of laser intensity increase since 1960	48
3.2	Schematic of the chirped pulse amplification method	51
3.3	Schematic of positive and negative chirp grating assemblies	52
3.4	Schematic of the ICLC Glass laser commercial oscillator	52
3.5	Schematic of ICLC Glass laser.	54
3.6	Single shot autocorellation pulse duration measurement.	55
3.7	Focal spot measurement.	56
3.8	Schematic of optical parametric amplification.	57
3.9	Schematic of Vulcan Target Area West beam-line.	58
3.10	Microscope capture of CR-39 tracks.	60
3.11	Schematic of HD-810 radiochromic film.	61
3.12	Results of RCF calibration against energy.	62

3.13 Results of RCF calibration against fluence.	63
3.14 Typical RCF stack with proton energy loss.	64
3.15 Typical RCF stack with proton ranges.	65
3.16 Results of RCF stack calibration.	66
3.17 Results of magnetic spectrometer calibration.	67
3.18 Schematic of dark-field schlieren imaging.	70
3.19 Schematic of a streak camera.	71
3.20 Schematic of a Mach-Zehnder interferometer.	73
3.21 Schematic of a Michelson interferometer.	73
3.22 Recovering phase information from interferometry.	74
3.23 Schematic of modified cryogenic gas nozzle.	76
4.1 Schematic of WDM Vulcan TAW experiment 2010	80
4.2 Data and NYM simulation of wire expansion	82
4.3 Measured pressure isentropes and Foord simulation for aluminium	84
4.4 Pressure-density plot from radiography data	85
4.5 Streaked XUV data from 2010 experiment.	86
4.6 Proton spectrum and energy deposition in Al from RCF stack.	88
4.7 Proton spectrum and energy deposition in Al from magnetic spectrometer measurements.	88
4.8 Proton and heavy ion spectrum and energy deposition in Al from magnetic spectrometer measurements.	89
4.9 Scan of Thomson parabola data	90
4.10 Raw streaked UV self-emission data.	91
4.11 UV emission data isolated from proton source.	92
4.12 Heating uniformity lineouts from UV emission.	92
4.13 Ion time of flight spread	94
4.14 Graphical representation of the numerical solution to the 1D heat transport equation.	95
4.15 2D plot of calculated deposited energy in Al.	97
4.16 Ion spectrum and energy deposition after ion filter.	97

5.1	Schematic of dual polarisation probing and duel heating beam setup.	103
5.2	Schematic of two perpendicular blast wave experiment.	104
5.3	Schematic of imaging setup for colliding blast waves experiment.	105
5.4	Raw Schlieren, interferometry and shadowgraphy images of a vertical blast-wave.	106
5.5	Raw images of horizontal beam delay scan.	106
5.6	Graph of horizontal beam delay against left separation.	107
5.7	Interferometry and Schlieren images of pre-heated blast wave.	108
5.8	Analysis of pre-heated interferometry image.	108
5.9	Graph of disassociation front velocity against time.	110
5.10	Nanoplasma model result for maximum ion energy from laser heated clusters.	111
5.11	SRIM model to fit the stopping power of argon gas.	112
5.12	Blast wave proton probing schematic.	116
5.13	Blast wave proton probing results.	116
5.14	SRIM output of proton path divergence through gas jet.	117
5.15	Vulcan 2007 blast wave experimental setup.	119
5.16	Absorption energy and length data.	120
5.17	Schlieren image of ionisation channels.	121
5.18	Schematic and data of a supercritical shock.	123
5.19	Shock surface structure in a Schlieren image.	124
6.1	Low energy proton spectra for various laser energies.	129
6.2	Schematic of a radial wire array experiment on MAGPIE.	130
6.3	MAGPIE proton pinhole and optical Schlieren images.	131
6.4	Schematic of virtual source measurement.	132
6.5	Proton energy dependence of angular divergence.	134
6.6	Schematic and results of low energy emittance measurement.	136
6.7	Experimental data SRIM uses for aluminium stopping power.	139
6.8	Schematic showing the useful range for stopping power determination of 6 μm aluminium.	140
6.9	Schematic for mono-energetic laser driven proton probing beam experiment.	143

7.1	Interferometric analysis of a 10 J shot into krypton clusters showing fringe reversal.	149
7.2	Hydrodynamic simulation of a collapsing conical blast wave.	151
7.3	Schlieren images from preliminary annular blast wave experiment.	152
7.4	Schematic of Cerberus laser system.	154
A.1	Normalised error of the magnetic spectrometer arising from finite slit width.	157

List of Tables

3.1	Several non-linear effects in dielectric materials arising from strong electric fields	50
4.1	Samples heated to warm dense matter state	81
5.1	Experimental gas jet parameters for similar conditions in different gases. . .	104

Chapter 1

Introduction

Recent rapid advances in laser technology have provided a means of achieving high energy density states in the laboratory through either direct laser heating or secondary radiation heating which were previously only accessible in the cores of planets, supernova or nuclear explosions. These smaller scale laboratory experiments allow a means to repeatably and reliably explore the parameter space of plasmas under extreme conditions on a relatively quick time scale. Comparisons can thus be made with these larger scale systems under certain hydrodynamically equivalent conditions [1] [2]. This thesis explores two distinct regimes where the parameter space is not well understood - warm dense matter and high energy density blast waves. While very different in terms of density and temperature the tools needed to create and probe these systems in the laboratory provides a unifying theme to the work presented here.

Warm dense matter typically exists as a transitory state in Earth based experiments and it is difficult to deliver the required heating on a fast enough time scale to also achieve the dense condition. We have developed a method to access this regime via laser generated ion heating, which requires less expense and reaches higher temperature than the typical method for achieving this via shock compression [3]. A growth in research in this regime is motivated by the inertial confinement fusion project at the National Ignition Facility [4] where understanding of warm dense matter is vital since all potential targets must pass through the warm dense matter regime on the way to ignition.

High energy densities are accessible in a university based laboratory through the unique ability of clustered gas to absorb higher than ninety percent of laser energy. Using the blast waves created by laser absorption in a cluster gas we can model phenomena such as supernova remnants [5] and the instabilities predicted and observed therein. This

thesis seeks to expand our characterisation of the blast wave system through experimental observation and the development and exploitation of new diagnostic techniques.

1.1 Contributors and Publications

All the experimental work carried out on the Vulcan Laser, a national facility, was supported by Central Laser Facility staff - both in maintaining the laser and providing technical support. Similarly the work on the Helen Laser required the aid of a large team of staff at AWE to run the laser and provide engineering support. The Imperial College Glass CPA Laser for the duration of the work described in this thesis was run, upgraded and maintained by the author. The Surrey Ion Beam Centre's Tandatron accelerator was maintained by C. Jeynes and the experimental concept with respect to the accelerator was developed with the aid of A. Clough. The magnetic spectrometer used in both detailed characterisation and proton spectroscopy studies was built at AWE following a design by E. Clark.

Low energy proton calibration of the spectrometer and radiochromic film was proposed and paid for by AWE, particularly W. J. Garbett and E. T. Gumbrell. The experimental work at Imperial was led by the author with support from S. I. Olsson Robbie. The analysis was performed in full by the author. At the time of writing a paper on the different methods of spectrometer calibration and response of radiochromic film to low energy protons is in preparation for submission to *Review of Scientific Instruments*.

The warm dense matter experiments described in chapter 4 were designed and run by AWE. The experiments were carried out by E. T. Gumbrell, S. F. James, T. M. Guymer, M. P. Hill, C. R. D. Brown, J. Fyrth, D. J. Hoarty and the author. The author was specifically responsible for the setup and analysis of the proton spectra data. At the time of writing, a paper on the method used for generating and diagnosing warm dense matter has been submitted to *High Energy Density Plasmas*.

The perpendicular blast wave experiments described in chapter 5 carried out on the Imperial College Glass Laser were conceived of and analysed by the author. The experiments were carried out with the aid of S. I. Olsson Robbie. An article describing the perpendicular blast wave method and observed ion wave was submitted to *High Energy Density Physics* in August 2011. This work was also presented at the 52nd Annual Meeting of the American Physical Societies Division of Plasma Physics by the author and also at an invited talk by R. A. Smith at the 3rd International Conference on High Energy

Density Physics.

The high energy scaling cluster blast waves experiments at the Rutherford Appleton Laboratories conceived by R. A. Smith were mainly organised and analysed by M. Hohenberger. The experiment was also carried out by J. Lazarus, R. E. Carley, D. R. Symes, A. S. Moore, E. T. Gumbrell, R. A. Smith and the author. A paper describing the first observation of a velocity instability was published in *Physics Review Letters* in November 2010 [6]. A more generalised paper on laser-driven radiative blast waves in clustered gas, which features some of this experimental work, was published in December 2009 in *High Energy Density Plasmas* [7].

A more theoretical blast wave paper, containing some of the work detailed in chapter 2, driven by modelling by D. R. Symes, R. Rodriguez and J. M. Gil was published in *Nuclear Instruments and Methods in Physics Research A* in December 2010 [8].

The Imperial College MAGPIE pulsed power generator work was carried out with the aid of the Plasma Physics Group, namely S. V. Lebedev and F. Suzuki-Vidal who were responsible for carrying out the MAGPIE wire implosion. Also, R. A. Smith for the proton pinhole camera concept which was analysed and deployed by S. Patankar and the author. The first ion emission spectra measurement of MAGPIE was carried out and analysed by the author. Planning is in progress for a temporally resolved ion energy spectra for publication.

Chapter 2

Theoretical Background

The aim of this chapter is to introduce the basic theoretical background and motivation for the experimental results presented in later parts of this thesis. The chapter starts by introducing the concept of warm dense matter and the current theoretical difficulties in accurately modelling it. This motivates the experimental work aimed at determining the equation of state in this regime in order to benchmark theoretical simulations. The basic definition of a plasma and related models are also briefly described in order to give some grounding for the scaling of laboratory produced radiative high energy density plasma blast waves to the astrophysical analogue described at the end of the chapter.

The middle of the chapter describes the physics behind the choice of target media for experimental applications. The absorption, heating and acceleration mechanisms for electrons and ions in overdense plasmas and isolated atomic clusters are also covered.

2.1 Beyond Condensed Matter

Beyond the better known solid, liquid and gas condensed phases exists warm dense matter (WDM) and plasma. Fig. 2.1 shows typical density and temperature values for these phases. It can be seen that the plasma state defines a broad stroke of temperature and density space - we are particularly interested in high energy density plasma for this thesis defined here in terms of energy per unit volume to be $\geq 10^{11} \text{ Jm}^{-3}$ [9].

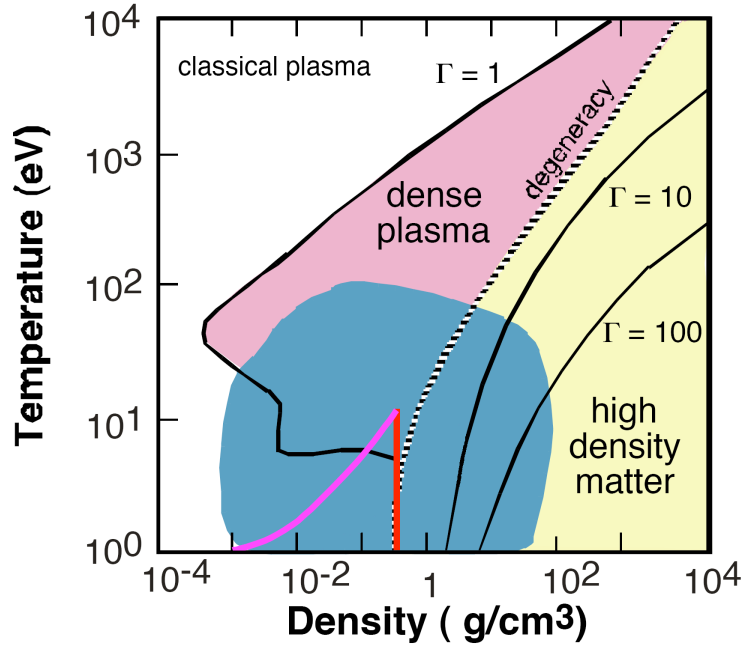


Fig. 2.1: Aluminium temperature-density phase diagram with the warm dense matter regime shown in blue. Regimes of the coupling parameter are shown and the yellow region marks where degeneracy becomes important in models. A sample heating isochore is shown in red along with a potential release isentrope in pink. Adapted from [10].

2.1.1 Definition of Warm Dense Matter

Warm dense matter can be broadly defined for the experimentalist as a material existing between 0.1 times solid density and ten times this value and at temperatures from 0.1 eV to 100 eV [11]. Typically it is viewed as a boundary state between condensed matter and plasma physics, as neither of these models accurately simulate warm dense matter. It is too energetic to be described by *ab initio* molecular dynamic calculations and too dense for standard plasma methods such as density functional theory to be used successfully [12].

More formally warm dense matter is a phase which is strongly coupled and mostly

degenerate. The strong coupling parameter Γ compares the potential energy between ions, V_{ii} , to the electrons' kinetic energy, E_{kin} .

$$\Gamma = \frac{V_{ii}}{E_{kin}} = \frac{Z^2 e^2}{r_0 E_{kin}} \quad (2.1)$$

where Z is the atomic number, e the elementary charge and the interaction radius, r_0 , is proportional to $\rho^{-1/3}$.

If $\Gamma < 1$, such as in a weakly coupled plasma, the plasma can be easily modelled as separate point charges. However, as coupling becomes stronger with decreasing E_{kin} and $\Gamma > 1$ particles interact, shifting their energy levels and ionisation potentials ensuring that modelling is non-trivial.

The partially degenerate nature of warm dense matter ($T \leq T_f$) results in non ideal behaviour where the pressure is no longer linearly dependant on temperature and the quantum mechanical equation of state for a Fermi gas is required, as the pressure is derived from the Pauli principle. Again this property makes modelling of the high kinetic energies and low collision rates in degenerate matter non-trivial.

2.1.2 The Equation of State for Warm Dense Matter

The difficulties described in §2.1.1 require current theoretical models to be pushed to their limits. As such, experimental data is needed to benchmark these strongly correlated, mostly degenerate many-body systems. Research into inertial confinement fusion [11], laboratory astrophysics [13] and geophysics [14] aims to directly measure the equation of state as well as optical and thermodynamic transport properties. This thesis concentrates in part on the poorly diagnosed warm dense matter equation of state problem [15].

Early equation of state tables [16] for the warm dense matter region interpolate between condensed matter and plasma models, as shown for aluminium in Fig. 2.2, resulting in significant inconsistencies internally and with respect to experimental data. To create a coherent equation of state for the whole spread of parameter space the Thomas-Fermi quotidian equation of state is typically used - this ignores much of the specifics of warm dense matter [15]. Recent interest has prompted research into models specifically designed to incorporate WDM physics such as quantum molecular dynamics [17].

Experimental benchmarking can be achieved through isochoric heating of a sample to the warm dense matter regime then observing the release isentrope, as shown in Fig. 2.1,

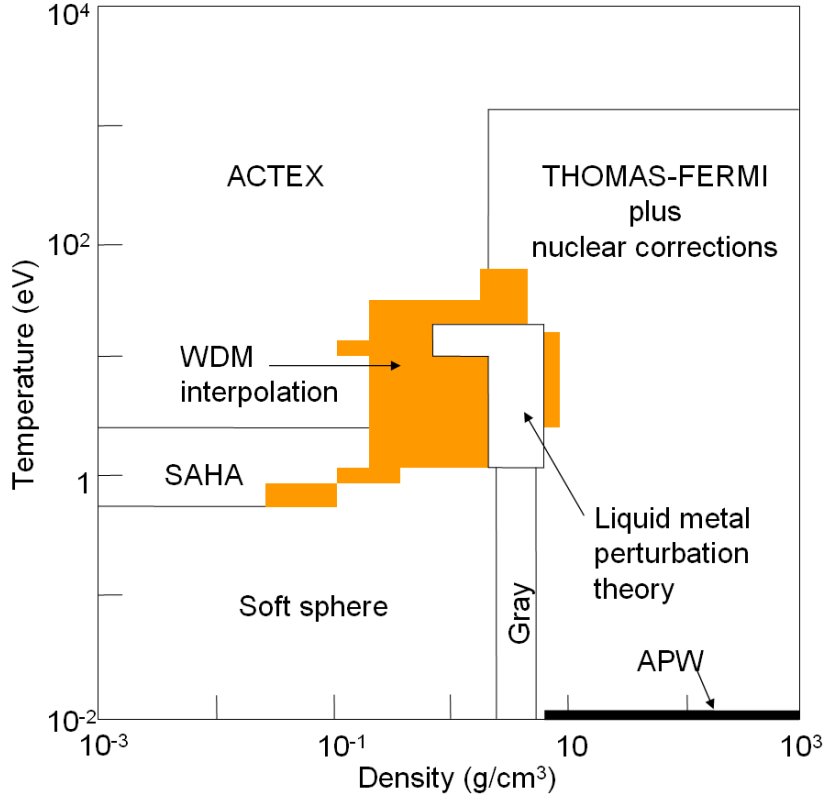


Fig. 2.2: Aluminium temperature-density phase diagram showing the equation of state models used to form the SESAME tables and the area where interpolation is used. Adapted from Holian *et al.* where more information on the ACTEX, SAHA, Gray, APW, Soft sphere and Thomas-Fermi equation of state models can be found [18].

with measurements of temperature and rate of expansion. *Isochoric* heating implies that no change in sample volume occurs during the heating process. Observing the release *isentropes* (constant entropy) implies that no energy is gained or lost by the sample after the heating process.

2.1.3 Methods for obtaining warm dense matter in the laboratory

Accessing the warm dense matter regime can be achieved through several methods, and two key requirements for well defined experimental investigations are good spatial uniformity and isochoric heating. Uniform heating requires energy deposition to penetrate through the bulk density of the material so that a single temperature state is achieved throughout the sample volume. Isochoric heating requires total energy deposition to occur before target hydrodynamic expansion, a time approximately equal to d/c_s , where d is the target thickness and c_s the heated sample speed of sound - of order a picosecond for typical micron scale targets [19].

Direct heating with ultrafast lasers [20] is within the required time-scale but is limited by the critical density surface, n_c , beyond which they will not penetrate and heat. For common laser wavelengths in the near-IR and visible this is less than 50 nm and would require ultra-thin targets where impurities and imaging resolution have a larger impact. An exception to this are x-ray free electron lasers which are currently being used to create WDM [10] however, these are relatively expensive and, at time of writing, are typically used in conjunction with a high power laser for probing the equation of state.

Shock heating [14] is the most common method for obtaining WDM. It can currently only achieve low temperatures (~ 5 eV) and is limited to on-Hugoniot measurements. A Hugoniot measurement is one which falls on the pressure-density curve defined by the Rankine-Hugoniot jump conditions and every material has a unique curve. More information on the Rankine-Hugoniot jump conditions can be found in Zel'dovich *et al.* [21]. Isentropic compression methods typically obtain few Mbar pressures [3]. Here the sample is compressed, removing the requirement for rapid heating before expansion.

Indirect laser heating through the generation and use of ions, electrons or x-rays has been made possible by increases in achievable laser power. These secondary sources of radiation have a duration on the order of the driving laser pulse (typically ≤ 1 ps) and can penetrate further into solid densities than standard laser wavelength photons. Protons in particular have relatively long absorption lengths and deliver near uniform heating, allowing a narrow range of density and temperature initial states to be created. Ions heat more efficiently than laser driven electrons or x-rays [22] - previous work has shown temperature of 20 eV are achievable [23].

In this thesis proton heating is used to access warm dense matter conditions, and also as a probe of electric and magnetic fields in laser generated shocks.

2.1.4 Definition of a plasma

Increasing temperature and density beyond the warm dense matter regime we reach the plasma state. Plasmas makes up 99% of visible matter in the universe (conveniently classing the recent discovery of dark matter [24] as non-visible) and as a result of this, and because of their many and varied technological applications is an extensive and diverse research area.

A plasma is a bath of electrons and ions that interact via the electromagnetic force. As these particles group together under their local electric fields their separation, r_0 , will

be limited by thermal kinetic motion. This gives rise to a fundamental plasma parameter known as the *Debye length*, λ_D , over which electric fields are shielded and particles can move defined as [9]:

$$\lambda_D = \left(\frac{e^2}{\epsilon_0 k_B} \left(\frac{n_e}{T_e} + \sum_i \frac{n_i Z_i^2}{T_i} \right) \right)^{-1/2} \quad (2.2)$$

where n_e and n_i are the electron and ion density respectively, T their temperature in eV and Z_i the charge state of ion species i . If we assume the ion mobility is negligible compared to the timescales of our interactions, usually the case in non-degenerate plasmas, we can simplify Eq. (2.2) by neglecting the second term to achieve the electron Debye length, λ_{De} which is used as the Debye length for this thesis:

$$\lambda_{De} = \sqrt{\frac{\epsilon_0 k_B T_e}{n_e e^2}}. \quad (2.3)$$

The Debye length gives rise to two properties that define a plasma - collective effects and quasi-neutrality. The volume of a Debye sphere with radius λ_D is charge neutral, or quasi-neutral, if $n_e = Z_i n_i$. This effect means particles move very differently than they would in vacuum. This is simply demonstrated in the case of a quasi-neutral two species plasma at equilibrium temperature T with a charge potential ϕ distributed according to Boltzmann statistics (proportional to $\exp(q\phi/k_B T)$ where q is charge). Here the charge density, ρ_q , around a test ion at the origin is:

$$\rho_q = \rho_z - e n_e \exp\left(\frac{e\phi}{k_B T_e}\right) + n_i Z e \exp\left(\frac{-Z e \phi}{k_B T_i}\right) \quad (2.4)$$

where the term ρ_z , the charge density of the ion at the origin, involves a delta function which integrates to a value of Ze . We can simplify this by using the condition of thermal motion being much greater than that due to electric force, $k_B T \gg q\phi$, known as the *weak coupling condition* to $\rho_q = Z + \frac{\phi \epsilon_0}{\lambda_D^2}$. Poisson's equation, $\nabla^2 \phi = \rho_q / \epsilon_0$, allows us to solve for ϕ at a distance r from the charge:

$$\phi = \frac{Ze}{4\pi\epsilon_0 r} \exp(-r/\lambda_D). \quad (2.5)$$

We immediately see one of the unique and defining characteristics of plasma that a charge's potential falls away exponentially quicker than that of an isolated charge in vacuum which scales as $1/r$. This *Debye shielding* has a requirement that leads to another fundamental value, the *plasma parameter* - the number of electrons in a Debye sphere,

$N_D = \frac{4}{3}\pi n_e \lambda_D^3$. Debye shielding requires $N_D \gg 1$ or $\lambda_D \gg n_e^{-\frac{1}{3}}$ - this is identical to the weak coupling condition [9]. As N_D increases towards infinity the plasma behaves more like an *ideal plasma* - analogous to an ideal gas.

The plasma parameters outlined above can be used to define an ideal plasma phase which is weakly coupled and quasi-neutral such that:

$$n_e^{-\frac{1}{3}} \ll \lambda_D \ll L \quad (2.6)$$

where L is the scale length equal to $n_e/|\nabla n_e|$.

2.1.5 Modelling plasma motion

The *single particle description* models the individual charge carriers' motion in an external electromagnetic field:

$$m \frac{d\mathbf{v}}{dt} = q \left(\mathbf{E} + \frac{\mathbf{v}}{c} \times \mathbf{B} \right) \quad (2.7)$$

where q , m and $\mathbf{v}$ are the charge, mass and velocity of the particle moving in the electric field $\mathbf{E}$ and magnetic field $\mathbf{B}$. This use of the Lorentz force, while accurate, requires huge computational power to track many particle systems and quickly becomes impractical without further simplifying assumptions. It is however a powerful tool for modelling magnetic field induced effects present in high energy density plasmas.

One common simplification is to use a statistical distribution such as the Boltzmann equation. An evolution of this *kinetic description* takes the collisionless Boltzmann equation - assuming collisional contributions to the electromagnetic field happen within the Debye sphere and are thus shielded. This is known as the *Vlasov equation* which uses self-consistent electromagnetic fields to obtain a phase space distribution function, $f(\mathbf{x}, \mathbf{v}, t)$:

$$\frac{\partial f}{\partial t} + \mathbf{v} \cdot \frac{\partial f}{\partial \mathbf{x}} + \frac{q}{m} \left(\mathbf{E} + \frac{\mathbf{v} \times \mathbf{B}}{c} \right) \cdot \frac{\partial f}{\partial \mathbf{v}} = 0. \quad (2.8)$$

The Vlasov equation is useful for computational work but typically in this thesis only the macroscopic properties of a plasma are required, not the detailed kinetic description. To this end we can use the *Euler fluid equations* which can be derived by taking velocity moments of the Vlasov equation. These moments can be rearranged to give the classical form of the Euler equations which are the continuity, force and energy equations for the density, mean velocity and pressure of particles with charge q_j and mass m_j [25], respectively:

$$\frac{\partial n_j}{\partial t} + \frac{\partial}{\partial \mathbf{x}} \cdot (n_j \bar{\mathbf{v}}_j) = 0 \quad (2.9)$$

$$n_j \left(\frac{\partial \bar{\mathbf{v}}_j}{\partial t} + \bar{\mathbf{v}}_j \cdot \frac{\partial \bar{\mathbf{v}}_j}{\partial \mathbf{x}} \right) = \frac{n_j q_j}{m_j} \left(\mathbf{E} + \frac{\bar{\mathbf{v}}_j \times \mathbf{B}}{c} \right) - m_j^{-1} \frac{\partial p_j}{\partial \mathbf{x}} \quad (2.10)$$

$$\frac{\partial p_j}{\partial t} + \bar{\mathbf{v}}_j \cdot \frac{\partial p_j}{\partial \mathbf{x}} = -\gamma \frac{\partial p_j}{\partial \mathbf{x}} \cdot \bar{\mathbf{v}}_j \quad (2.11)$$

where n is the density, $\bar{\mathbf{v}}$ is the average velocity, $\mathbf{x}$ is the position vector, γ is the polytropic index and p is the isotropic pressure. Here we have assumed the pressure is isotropic and follows the adiabatic equation of state for an ideal gas, $p = n^\gamma$. The polytropic index can be defined as the ratio of specific heat at constant pressure to specific heat at constant volume for an adiabatic process. Using thermodynamics to express the specific heat ratio in terms of the number of degrees of freedom of particles, ϑ , we see $\gamma = \frac{\vartheta+2}{\vartheta}$ for a classical ideal gas where $T_i \ll T_e$ - not Bose-Einstein or Fermi gases. So for a classical, ionised, non-relativistic gas $\gamma = 5/3$.

The collisional coupling between ions and electrons reduces as temperature increases and density decreases, leading them to act as two separate fluids. Here the single fluid model given above breaks down and fails to properly simulate effects such as plasma oscillations [9]. To overcome this the *two fluid model* is used, the main framework the reader will encounter in this thesis. Here two phase space distribution functions are established, f_e and f_i , for the electrons and ions respectively, which have their own properties and act ‘independent’ of each other. The two-fluid equations as a multi-fluid model is necessary to accurately describe the interaction of an intense laser with a plasma [25] - the focus of this thesis.

Hence with equations (2.9), (2.10), (2.11) and the Maxwell equations to relate electromagnetic fields to the plasma charge and current densities we have a complete description of the two field model for a non-collisional thermal equilibrium plasma consisting of electrons and ions (in a single charge state). As the plasma becomes collisional the full Boltzmann equation must be used by replacing the zero term in Eq. (2.8) with $\left(\frac{\delta f_{ei}}{\delta t} \right)_{coll}$ which is the rate of change of each species phase space distribution function due to collisions.

One of the limitations of the Euler equations is that they are moments of each other. The 0th order moment will contain terms from the 1st, the 1st from the 2nd and so *ad infinitum*. To account for this we set the contribution to the 2nd order from the 3rd order to be zero. The 3rd order is heat flux - our system will not model heat flux. Although the Euler fluid equations will model most shock phenomena, occasionally we must use *general*

fluid equations which account for phenomena such as radiation flux and heat transport.

2.2 Laser-Matter Interactions

In this section we describe the processes behind coupling of laser energy to a target to initially ionise it in order to produce a plasma and then to further couple energy to achieve high energy density plasmas or induce ion acceleration via high temperature electrons. The choice of target media is also discussed with regards to the most efficient material for these applications within laser-driven plasma physics.

2.2.1 Ionisation Mechanisms

Laser pulses can ionise material via several mechanisms. At low laser intensities the *photoelectric effect* is dominant. Here an atom absorbs a photon of energy $E = h\nu$, where h is Planck's constant and ν is the light frequency, and in turn emits an electron with kinetic energy $E_{kin} = h(\nu - \nu_0)$. ν_0 is the minimum photon frequency necessary to elevate an electron from the atomic Coulomb potential well.

As the light intensity increases beyond $\sim 10^{10} \text{ Wcm}^{-2}$ *multiphoton ionisation* can occur. Here many photons are needed to elevate the electron from its Coulomb potential well. The electron undergoes virtual transitions of energy $h\nu$ until it has an energy greater than the ionisation energy. These virtual energy levels have a very short lifetime, and so this process can only occur at high intensity ($>10^{10} \text{ Wcm}^{-2}$). *Above threshold ionisation* can occur at higher intensities ($\geq 10^{14} \text{ Wcm}^{-2}$) whereby more photons then needed for MPI promote the electron into the continuum with additional kinetic energy. However, as the light intensity is increased further *optical field ionisation* can occur.

The Keldysh parameter compares the ionisation energy to the ponderomotive energy (the energy given to the electron by the laser electric field) and can be used as a quantified indicator of the dominant mechanism [26]:

$$\lambda_K = \sqrt{\frac{\epsilon_i}{2U_{pondero}}} \quad (2.12)$$

where ϵ_i is the ionisation energy and $U_{pondero}$ is the ponderomotive potential. For $\lambda_K \gg 1$ the dominant process is multiphoton ionisation (MPI).

For $\lambda_K \ll 1$ optical field ionisation can occur. Here, the potential due to the laser field is high enough to distort the Coulomb potential well, typically for intensities $>10^{14} \text{ Wcm}^{-2}$.

This distorted potential promotes quantum mechanical tunnelling known as *tunnel ionisation*. As the field increases even further the potential barrier can be suppressed to below the bound electron energy. The critical intensity for this so called *barrier suppression ionisation* is [27]:

$$I_{threshold} = \frac{\pi^2 c \epsilon_0^3 \epsilon_i^4}{2Z^2 e^6} \quad (2.13)$$

where Z is the charge state of the final ion. This means that barrier suppression ionisation is the dominant mechanism for most laser accelerated electrons in this thesis. The intensities required for these ionisation processes are several orders of magnitude lower than typical peak laser intensities employed in the experiments described in this thesis. This means that usually a low level laser pre-pulse will establish a surface plasma at a solid surface before the main pulse arrives. Figure 2.3 shows the physical process behind the above mechanisms.

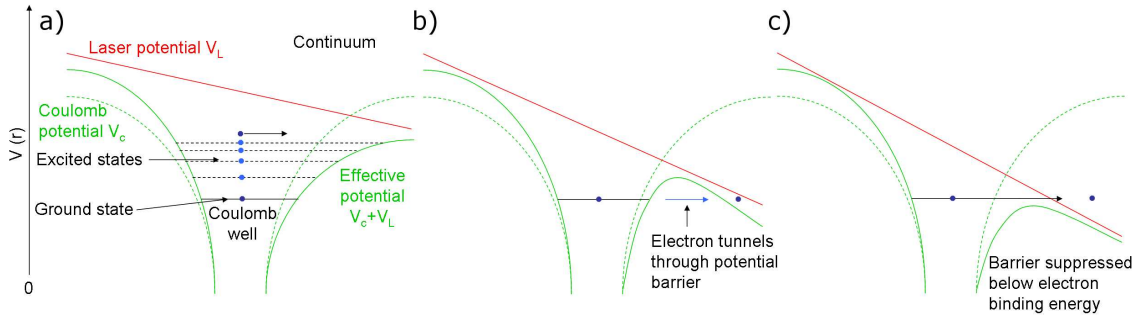


Fig. 2.3: Diagram showing optical ionisation processes which enable the electron to escape the potential well: a) multiphoton ionisation b) tunnelling ionisation c) barrier suppression ionisation.

The optical ionisation processes described above will ionise surface electrons from a solid target. However, as the laser propagates into the solid target a critical density, $n_c = \frac{\omega^2 \epsilon_0 m_e}{e^2}$, will quickly be reached, at which point reflection occurs and ionisation deeper into the target by the above processes cannot occur. For these dense materials we rely on *collisional ionisation*. Here the free electrons from optical field ionisation, or other processes, collide with atoms and ionise them as discussed in §2.2.4.

2.2.2 The Ponderomotive Force on an Electron in a Laser Field

The *single particle description* given in §2.1.5 can be used to explore the force on an electron in a plasma exerted by an external laser field described as $\mathbf{E}(\mathbf{r}, t) = \mathbf{E}_r \cos \omega_L t$,

$\mathbf{B}(\mathbf{r}, t) = \mathbf{B}_r \cos \omega_L t$ where ω_L is the angular frequency of the laser and $\mathbf{E}_r$ the spatially varying part of the electric field with $\mathbf{B}_r$ similarly that of the magnetic field. The non-relativistic Lorentz force gives the equation of motion as:

$$m_e \frac{\partial \mathbf{v}_e}{\partial t} = -e(\mathbf{E}_r \cos \omega_L t + \mathbf{v}_e \times \mathbf{B}_r \cos \omega_L t). \quad (2.14)$$

Using $c = \frac{\omega_L}{k}$ and Faraday's law to show that $\mathbf{E}_r = \mathbf{B}_r \frac{\omega_L}{k} = c\mathbf{B}_r$ we see that upon integration the magnetic field term is negligible for $v_e/c \ll 1$. This allows an expression for electron velocity:

$$\mathbf{v}_e(\mathbf{r}, t) = \frac{e\mathbf{E}_r}{m_e\omega_L} \sin \omega t. \quad (2.15)$$

Thus the electron will oscillate with a velocity, $v_{quiver} = \frac{eE_r}{m_e\omega_L}$, known as the *quiver velocity*. Substituting this velocity back into the non-relativistic equation of motion, multiplying through by m_e and taking the cycle average we obtain the non-relativistic ponderomotive force, $\mathbf{f}_p$ [26]:

$$\mathbf{f}_p = -\frac{e^2}{4m_e\omega_L^2} \nabla \mathbf{E}_r^2. \quad (2.16)$$

This force will push the electrons radially away from regions of high laser intensity, such as a focal spot, with velocities on the order of the quiver velocity. Using $U_p = \frac{1}{2}m_e v_{quiver}^2$ as a definition of the *ponderomotive potential* we can use practical measurable units to show the relative importance of the ponderomotive force:

$$\frac{U_p}{1 \text{ keV}} \approx \frac{I}{10^{16} \text{ Wcm}^{-2}} \left(\frac{\lambda^2}{1 \mu\text{m}} \right) \quad (2.17)$$

where I is the intensity of laser with wavelength λ . For intensities $\geq 10^{18} \text{ Wcm}^{-2}$ the electron quiver velocity approximates to the speed of light, the non-relativistic limit is not valid and we must include the magnetic term and changes to m_e in the Lorentz force. The quiver velocity becomes $v_{quiver} = \frac{eE_0}{\gamma m_e \omega}$ where $\gamma = 1/\sqrt{1 - v_e^2/c^2}$ is the Lorentz factor for an electron with total velocity v_e and energy $E_e = \gamma m_e c^2$. Kruer considers the equation of motion for a relativistic element in the fluid model to show that the longitudinal component of the momentum is given by [25]:

$$\frac{\partial \mathbf{p}}{\partial t} = e\nabla \Phi - m_e c^2 \nabla (\gamma - 1). \quad (2.18)$$

Here the first term is the charge separation force where Φ is the electrostatic potential and the second term is the ponderomotive force with an associated ponderomotive potential,

U_p , giving the energy an electron gains as it oscillates in the laser field:

$$U_p = (\langle \gamma \rangle - 1)m_e c^2. \quad (2.19)$$

The electric and magnetic components act perpendicular to each other and the laser propagation direction. This results in the non-relativistic one dimensional oscillation becoming a two dimensional figure of eight oscillation in the relativistic regime. Thus the electric term contributes a drift velocity component away from the high intensity focal spot and the magnetic term a component in the direction of the laser propagation [25]. This electron acceleration can act to create a channel depleted of electrons, a process known as *ponderomotive channelling*. The acceleration of an electron will continue until it is equal to the deceleration resulting from the attractive space charge force.

In order to show the laser intensities at which relativistic effects become important the relativistic v_{quiver} can be substituted into the Lorentz factor γ :

$$\gamma = \sqrt{1 + \frac{e^2}{2\pi^2\epsilon_0 m_e^2 c^5} I \lambda^2} = \sqrt{1 + a_0 I_0} \quad (2.20)$$

where I_0 is the laser intensity at which the quiver velocity is equal to the speed of light and a_0 is the normalised vector potential. For order of magnitude calculations we can use:

$$\gamma \approx \sqrt{1 + \frac{I \lambda^2}{1.4 \times 10^{18} \text{Wcm}^{-2}}} \quad (2.21)$$

to show that only for intensities greater than 10^{18}Wcm^{-2} do relativistic effects become important.

This process of directly coupling energy from the laser to the electrons in a plasma is only possible before the critical density surface is reached and light is reflected, these plasmas are termed underdense. Beyond this limit, where plasmas are known as overdense, additional mechanisms become dominant - these are discussed later in this chapter.

2.2.3 Different Target Media for Laser Interaction Experiments

The wide range and properties of available materials have all been exploited in various laser interaction experiments. The choice of target will directly affect energy coupling and release mechanisms. One of the main properties we consider is the target density and we choose this with respect to the application. Here we look at a typical range of materials,

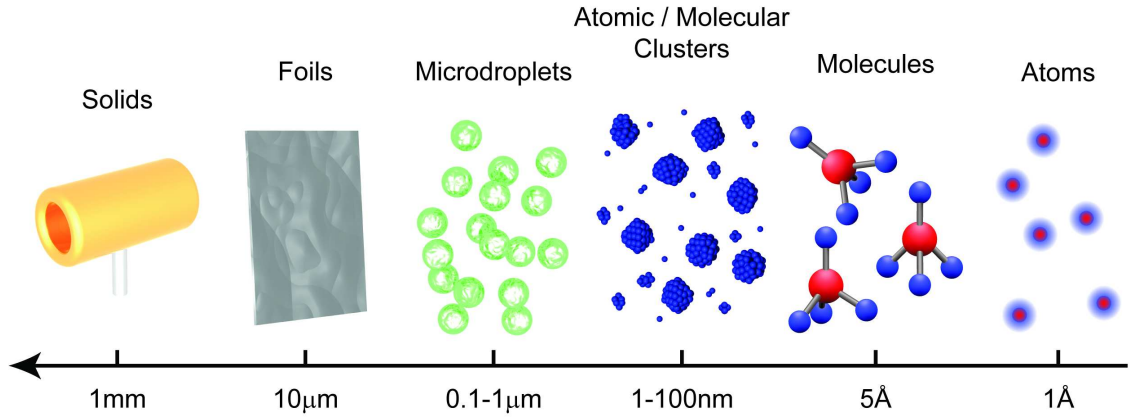


Fig. 2.4: Laser interaction target media from low density atomic gases to foils and solids. Diagram courtesy of Dr. M Hohenberger [28].

shown in Fig. 2.4 and discuss their applications.

Relatively low density atoms and molecules with a much smaller size than the laser wavelength absorb very little of the laser energy so are ill suited to the production of high energy density plasmas. However, their low density means laser radiation can penetrate through the target, and direct laser acceleration and heating occurs throughout the sample. This has given rise to an immense research field investigating nanometre scale laser-plasma accelerators with the potential of replacing expensive conventional kilometre scale linear accelerators. This culminated recently with experiments producing the first *GeV* electron beam from such a laser driven scheme [29]. Another important research field requires low density targets in which relatively low intensity lasers ($\leq 10^{14} \text{ Wcm}^{-2}$) are used to produce broadband high harmonics which can be potentially compressed to attosecond pulses [30].

Clusters consist of a few tens to many thousands of atoms or molecules bonded by Van der Waals forces or hydrogen bonding. They present a unique and unusual feature in that they have low average density but almost solid local densities at the cluster sites. Thus the laser will penetrate through the sample but efficient absorption ($\geq 90\%$) will still occur. These properties make them well suited to high energy density plasma creation and shock research, as such they feature extensively in this thesis. They have also been used as a source of x-rays [31] and for small scale laser driven fusion studies [32].

Heating of microdroplets occurs via surface laser absorption followed by fast heat conduction as their size is on the order of the laser wavelength. Local sub-wavelength field enhancements have been observed, causing emission of high energy ions [33] and x-ray production [34].

Laser energy can couple efficiently to surface plasmas on thin foils ($\sim 10 \mu m$) however beyond the critical surface further heating is reliant on potentially less efficient resonant effects. Microdroplets may seem a more efficient method of ion generation, however thin foils are the typical target of choice for their ease of use and will feature strongly in this thesis. The mechanism for ion production in thin foils can result in a high energy, low emittance beam which is more suitable for ion heating and imaging applications [35].

Bulk solids are used as targets to investigate inertial confinement fusion [36], shocks [37] and equation of states [38].

2.2.4 Energy Transfer Mechanisms in an Overdense Plasma

§2.2.2 describes direct coupling of laser energy to an electron via the Lorentz force. Neglecting the magnetic term and using Maxwell's equations we can derive the dispersion relation for the propagation of an electric wave through a plasma [25]:

$$\omega_L^2 = k^2 c^2 + \omega_p^2 \quad (2.22)$$

where $k = 2\pi/\lambda$ is the angular wave number and the non-relativistic electron plasma frequency is $\omega_p = \sqrt{\frac{n_e e^2}{\epsilon_0 m_e}}$. We can see that for direct laser acceleration to occur $\omega > \omega_p$. The critical surface where $\omega = \omega_p$ is known as the *critical density surface*, $n_{crit} = \frac{\epsilon_0 m_e \omega^2}{e^2}$. For $n_e > n_{crit}$, where the laser will not penetrate, the plasma is overdense.

In this overdense regime where direct laser acceleration cannot occur, various secondary mechanisms become dominant at different laser intensities, which are able to deliver heat through to the bulk of the material.

As an oscillating electron in the laser field interacts with a relatively massive ion the motion will be damped and it will lose some of its energy by collisional absorption, a process known as *inverse bremsstrahlung absorption*. Hence energy is transferred from the laser to the plasma. The non-relativistic spatial damping rate of the laser wave energy, κ_{coll} , is derived by *Wilks and Kruer* [39]:

$$\kappa_{coll} \propto \frac{Z n_e^2}{T_e^{3/2}} \left(1 - \frac{n_e}{n_{crit}} \right)^{-1/2}. \quad (2.23)$$

Thus, collisional absorption is greater for high atomic number, high density and low temperature plasmas. The fractional intensity absorption, f_{ib} , can be found from a simple

model [39]:

$$f_{ib} \approx 1 - \exp\left(\frac{32}{15} \frac{L}{c} \nu_{ei} n_{crit}\right) \quad (2.24)$$

where the electron-ion collision frequency is $\nu_{ei} \propto \frac{n_{crit} Z}{(k_B T_e)^{3/2}}$ and L is the plasma scale length defined as $L = \frac{n_e}{|\nabla n_e|}$. This electron damping term is the dominant plasma absorption mechanism in sub-critical densities up to intensities of $\sim 10^{15} \text{ Wcm}^{-2}$. Beyond this intensity the quiver velocity of the electrons approaches their thermal velocity and collisional absorption plateaus due to a decrease in the collision cross section.

Collisionless absorption processes can penetrate into overdense plasma regions within a solid target. The plasma dispersion relation for electromagnetic waves makes it clear that electromagnetic waves cannot penetrate beyond the critical density surface. In the case that k becomes imaginary, in the overdense regime, the wave will become evanescent. This evanescent wave can however drive *resonance absorption*. As the incident light wave is reflected near the critical density surface, a component of the electric field parallel to the density gradient can tunnel through to the critical density surface and excite plasmons that decay by particle trapping and collisions to heat electrons. The fraction of energy transferred by ‘classical’ resonance absorption for light incident at an angle θ to the critical surface is [26]:

$$f_{ra} = \frac{1}{2} \left(2.3 \tau e^{-\frac{2}{3} \tau^3} \right)^2 \quad (2.25)$$

where $\tau = \left(\frac{\omega L}{c}\right)^{\frac{1}{3}} \sin\theta$ and L is the plasma scale length. In order to maximise this energy coupling function $\tau \approx 0.8$. Thus an optimum angle can be found $\theta = \arcsin\left(\frac{0.8\lambda}{(2\pi)^{\frac{1}{3}} L}\right) \approx 25^\circ$ (for high intensities where L is of the order of the laser wavelength, λ).

As the plasma scale length decreases f_{ra} decreases. Thus for small L and short pulse lengths ($< 1 \text{ ps}$) where the plasma has very limited time to expand and the interaction takes place on a short, steep density gradient we rely on *vacuum heating*. Here the electrons in the sub-critical part of the density gradient are dragged into vacuum by the oscillating electric field of an obliquely incident laser. As the field changes direction the electrons can gain enough kinetic energy to penetrate the critical density surface, where they are no longer subject to this field as it reverses direction. Thus a fraction of laser energy is coupled into the plasma. For a p-polarised beam with intensities $\sim 10^{16} \text{ Wcm}^{-2}$ this fraction can be as much as 70% [27]:

$$f_{vh} = \left(\frac{\eta}{2\pi} \right) \left(\frac{v_{osc}^3}{v_l^2 c \cos \theta} \right) \quad (2.26)$$

where η is the efficiency constant, v_{osc} is the electron oscillation velocity and $v_l = eE/m_e\omega$. This mechanism saturates at high intensities where $v_l \geq c$ due to $\mathbf{v} \times \mathbf{B}$ deflections effects.

At relativistic intensities where the magnetic term of the Lorentz force cannot be neglected ($\geq 10^{18}$ Wcm $^{-2}$) mechanisms such as *relativistic $\mathbf{j} \times \mathbf{B}$ heating*, *hole boring* and *magnetic field generation* becoming increasingly dominant. $\mathbf{j} \times \mathbf{B}$ heating is particularly interesting as it is most effective at normal incidence to the plasma, unlike the other mechanisms described above which require obliquely incident light. The electrons are accelerated twice by the magnetic field in their figure of eight orbit under the Lorentz force per laser cycle. This can provide enough energy in the direction normal to the critical surface to accelerate the electron into the overdense plasma.

The processes described above are some of the more common mechanisms for electron acceleration through laser interactions. It should be noted however that this acceleration will not always be along the direction of laser propagation and can be heavily damped by the slower moving, heavier ions, or modified by large scale electric or magnetic fields created in the plasma, by for example bulk flow of electrons and ions.

2.2.5 Laser Driven Ion Acceleration from Solid Targets

Until such time as laser intensities of 10^{24} Wcm $^{-2}$ become available to directly accelerate ions through the Lorentz force we rely on accelerated electrons to in turn accelerate the ions through space charge separation. The electric fields generated by electron acceleration can be comparable to those within the laser field. The ion response is $\sqrt{m_i/Zm_e}$ slower due to their higher inertia.

It is generally accepted that ion acceleration can currently occur by two mechanisms in solid target interactions - *target normal sheath acceleration* (TNSA) and more recently *radiation pressure acceleration* (RPA). The ion energy of the TNSA process is dominated by T_e and n_e so maximum proton energy scales as $(I\lambda^2)^{1/2}$ where as RPA based on momentum conservation scales as $\alpha(I\tau/\sigma)^\alpha$ with τ the pulse duration, σ the areal mass of foil and α equal to 2 (or 1/3 in the ultrarelativistic regime) [40]. The scaling would suggest RPA is the more desirable mechanism. However, RPA which uses the same mechanisms as a light sail to transfer momentum from the laser to the target requires nanometre thickness foils and very high contrast pulses to avoid destroying the foil before the main

pulse arrives to accelerate it as a slab. The resulting energy spectrum from RPA can be quasi-monoenergetic compared to TNSA. For this thesis TNSA, the most commonly used mechanism [41] [42], is the mechanism of choice to create proton beams used to heat material to the warm dense matter state with a wider energy spectrum to increase uniformity of heating. Also the lasers used in this thesis do not have the capabilities to access the RPA regime.

TNSA relies on electric fields strong enough to create a plasma on the rear target surface and to accelerate electrons out of the target into vacuum [43]. The rear surface plasma ionises surface contaminants - particularly hydrogen. The space charge separation establishes electric fields, $E \approx \frac{k_B T_{hot}}{e \lambda_D}$, typically with values of 10^{12} Vm^{-1} which can then produce MeV ion energies. Although hydrogen ions are preferentially accelerated due to the low mass to charge ratio, other ions such as carbon and oxygen are also found in the resulting ion beam. This process can occur on both the front and the rear surface but recent work removing rear surface contaminants has shown that 99% of the ion beam comes from the rear surface [44]. Experiments have shown the ions are accelerated normal to the surface by using a wedged target and observing acceleration normal to both surfaces [41]. A model by Mora [42] which performs an analysis of the evolution of the electric field enables us to estimate a maximum ion velocity if we assume the acceleration time, t , is the same as the laser pulse length [45]:

$$v_{max} \approx \sqrt{\frac{4Zk_b T_e (\ln(\tau + \sqrt{\tau^2 + 1}))^2}{m_i}} \quad (2.27)$$

where $\tau = \omega_{pi} t / \sqrt{2e^1}$ for ions with charge Z feeling a plasma ion wave of frequency ω_{pi} for the duration, t , of the laser pulse. T_e represents the temperature of the fast electrons which accelerate the ions and is given by the laser ponderomotive potential [46]:

$$T_e = m_e c^2 \left(\sqrt{1 + I \lambda_{\mu m}^2 / 1.37 \times 10^{18}} - 1 \right). \quad (2.28)$$

Mora's self-similar model also gives the number of accelerated ions, N , per unit energy [42]:

$$\frac{dN}{dE_i} = \frac{n_0 t}{\sqrt{2E_i}} \exp \left(-\sqrt{\frac{2E_i}{Zk_b T_e}} \right) \quad (2.29)$$

where E_i is the ion energy and n_0 the initial surface plasma ion density. Mora predicts a sharp cutoff to the spectrum at the maximum ion energy as a result of a leading 'ion front'

formed by the sheath field which remains well defined due to a shielding effect preventing further acceleration [47].

Figure 2.5 shows the mechanisms for ion acceleration described above. It also shows that we must consider how ion transport through the target occurs. As the fast electrons travel through the target they create large magnetic fields which can focus them, in the forward direction [48]. These same magnetic fields however will deflect the ions. This deflection is used to explain ring structures commonly observed in proton beams as shown in figure 2.6 [49]. Though it should be noted that recent literature suggests other explanations for these rings - namely that the feature mostly vanishes when taking into account higher energy proton contributions or is caused by other effects such as the time dependant structure of the accelerating field [50] [51].

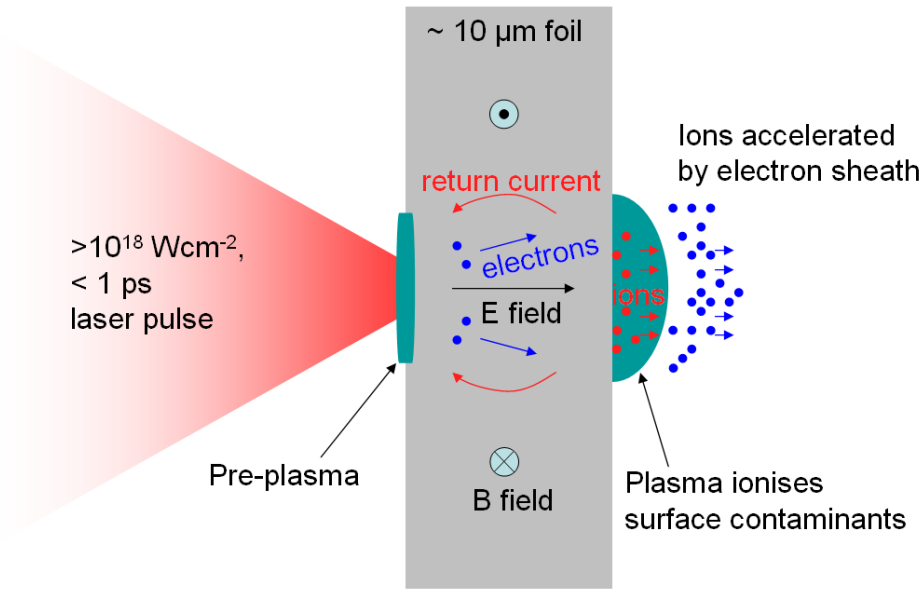


Fig. 2.5: Simplified diagram of the key mechanisms occurring in a laser interaction with a solid target leading to forward ion acceleration. The return current and magnetic fields established by the laser driven electron transport affect the ion beam quality and are highly dependent on foil thickness.

TNSA ion beams have typical maximum proton energies of 60 MeV with an approximately exponential spectral distribution peaked at low energies [41]. Depending on target and laser parameters the ion beam typically contains $\sim 15\%$ heavier ions in addition to protons. The beam typically has a divergence of $\sim 20^\circ$ with the higher energies confined to the beam centre [53]. Acceleration happens on time scales similar to the laser pulse, $\sim 1 \text{ ps}$, however the wide energy spectrum rapidly results in time of flight spreading [54].

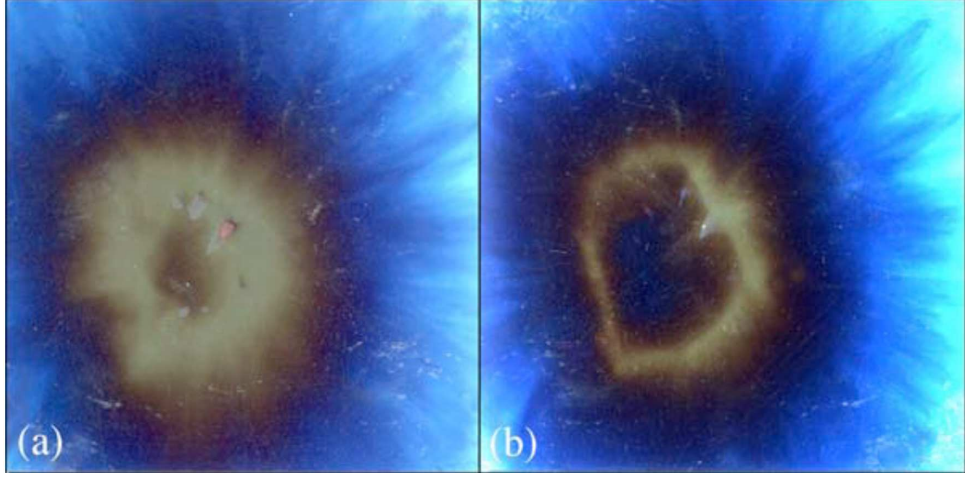


Fig. 2.6: Radiochromic film exposed to protons forward accelerated from a laser foil interaction performed by Clark [52]. a) Shows dose due to protons of energy greater than 1 MeV whereas b) placed behind it shows dose due to protons of energy greater than 4 MeV. The ring shown in the saturated copper colour can clearly be seen, this is attributed to magnetic deflection within the target.

Transverse emittance of less than 0.004 mm mrad for proton energies greater than 10 MeV have been demonstrated which is an improvement of more than hundred times over conventional accelerators [55]. The conversion efficiency has been previously measured as $\sim 12\%$ [41] [46].

2.2.6 Electron and Ion Transport and Stopping in Matter

For both TNSA, where the accelerated electrons propagate through the target foil, and in proton heating to warm dense matter it is vital to know how the particles are stopped in matter in order to determine the energy deposited and their range in material. Blast wave transport through matter and the affect of pre heated stopping power is also not well understood.

Stopping power is governed by the energy loss of a particle per unit distance and depends on the incident energy and material properties. For a given material it scales linearly with density.

At keV energies and above the particles suffer mainly small angle collisions and very few large angle collisions ($\ll 1\%$ in electrons and $< 10\%$ in ions). Hence we can simply model the propagation as an essentially one dimensional *continuous slow down approximation* where the range of a particle is given by [19]:

$$x_{CSDA} = \frac{\int_{E_0}^{E_{min}} S(E)^{-1} dE}{\rho} \quad (2.30)$$

where $S(E)$ is the stopping power of the material for a given energy E and E_{min} is the minimum energy at which the particle is considered absorbed. The stopping power is semi-empirical and huge databases exist where ranges and stopping powers for cold materials can be looked up - for electrons by NIST [56] and for ions SRIM [57]. However, in plasmas and warm dense matter very little empirical data exists and we investigate methods for determining warm stopping powers in this thesis.

2.2.7 Laser-Cluster Interaction at High Intensities

The unique high absorption ($> 90\%$) property of clusters for sub-picosecond laser energy and the resulting high energy electrons, ions and radiation is non-trivial to model. Ditmire *et al.* created the *nanoplasma model* which has been remarkably successful at simulating most observations given its many simplifying assumptions [58]. Here the clusters are considered as spheres of uniform density and radius R_c where the model is valid for $R_c \gg \lambda_D$. However the model also assumes constant temperature across the sphere so is most relevant for $R_c < \text{laser skin depth}$ which is the depth the laser penetrates into a plasma (c/ω_p). The energy is coupled into the sphere as if it was an overdense plasma - through optical ionisation and collisional inverse Bremsstrahlung heating.

As $R_c < \lambda$ the field inside the sphere can be described as that of a dielectric inside a constant field, E_0 [59]:

$$E = \frac{3E_0}{|\epsilon + 2|} \quad (2.31)$$

where the plasma dielectric constant, $\epsilon = 1 - \frac{\omega_p^2}{\omega(\omega + i\nu_{ei})}$ with ν_{ei} the electron-ion collision frequency, is derived from the Drude model which describes the plasma as an electron fluid damped by collisions with ions.

Eq.(2.31) shows that the heating from the electric field will be maximum when $|\epsilon + 2| = 0$. This resonant condition is achieved for the cluster's *Mie frequency*, $\omega_p = \omega\sqrt{3}$. We can state this more practically as the condition $n_e/n_{crit} = 3$. Fig. 2.7 shows this resonance inside the cluster and also a shielding effect as $n_e/n_{crit} > 6$.

The cluster will initially expand either through hydrodynamic pressure, p_{hydro} , due to the internal hot electron temperature or through Coulomb pressure, $p_{Coulomb}$, due to fast electrons escaping and the resulting charge buildup causing a *Coulomb explosion*. Charge buildup can be modelled by integrating the Maxwellian free streaming rate for electrons with velocities large enough to escape the cluster potential over the cluster surface. The

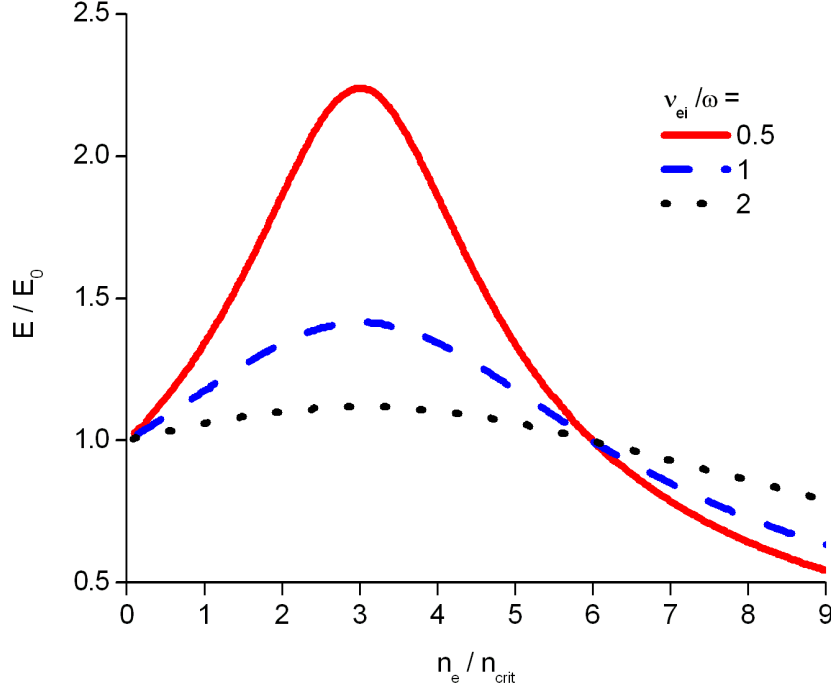


Fig. 2.7: Nanoplasma model output showing a transient resonance in the electric field within a cluster at $n_e/n_{crit} = 3$ as the density increases followed by a shielding effect as $n_e/n_{crit} > 6$.

hydrodynamic pressure is $p_{hydro} = n_e k_B T_e$ and scales as $1/R_c^3$. The Coulomb pressure is derived by treating the cluster as a conductor with charge Qe over the surface leading to a stored energy density per unit area of surface, $p_{coulomb} = \frac{1}{2} \frac{1}{(4\pi)^2 \epsilon_0} \frac{Q^2 e^2}{R_c^4}$, which scales as $1/R_c^4$. The hydrodynamic term typically dominates for larger radii.

Combining the pressure terms, $p = p_{hydro} + p_{Coulomb}$, we can obtain the equation of motion for the cluster, $4\pi R_c^2 p = \frac{4}{3} \pi R_c^3 n_i m_i \frac{\partial^2 R_c}{\partial t^2}$ which leads to [27]:

$$\frac{\partial^2 R_c}{\partial t^2} = \frac{3}{R_c} \frac{p}{n_i m_i}. \quad (2.32)$$

The nanoplasma model uses the electric field heating and expansion radius derived above, along with calculated electron cooling due to expansion and ion collision. An example nanoplasma model output from Dr Hohenberger is shown in Fig. 2.7 [28]. As the cluster expands and the density drops the $n_e/n_{crit} = 3$ condition will be met again and a second resonance is modelled.

An advanced version of the nanoplasma model does not assume a uniform density and allows for an angle dependant electric field [60]. This purely hydrodynamic model still sees an initial resonance at $n_e/n_{crit} = 3$ but has a smaller amplitude ($\sim 1 : 100$) that lasts ~ 50 times longer due to the non-uniform electric field producing a radial

ponderomotive force which compresses the cluster core to above critical density for longer. The second resonance is also not simulated. Both models agree on $> 90\%$ absorption efficiencies however the latter model can simulate observed phenomena such as cluster self-focusing where the susceptibility, χ , is reversed [61]. More advanced models are necessary to accurately model ultrafast few femtosecond laser interactions and angular dependent phenomena [62] [63].

2.3 Shocks in Plasmas

In this section the fluid description of a plasma is carried forward to describe shocks and blast waves produced in laser-plasma interactions. The concept of scaling to astrophysical regimes and the necessary conditions are also described. We then highlight the astrophysical regime we are able to model experimentally with blast waves from a clustered gas.

2.3.1 Shock Waves and Blast Waves

When energy is abruptly deposited in a fluid, such as occurs in the interaction of a high intensity picosecond laser pulse with an extended gas, a shock is often formed. Many astrophysical events, such as supernova, also create shocks [64] [65]. A shock is a sudden transition in the pressure, temperature and density properties of a fluid that arises when a fluid is compressed as it is forced to move faster than the adjacent ambient fluid can flow out of the way, faster than the speed of sound in the pre-shocked material, c_{s1} . As the shock wave propagates it will transport energy, heating and potentially ionising matter. If the initial energy deposition is rapid, as in the case of a short laser pulse, the initial compression wave will be followed by a rarefaction caused by the expansion of the fluid [9]. The rarefaction wave moves at the post-shocked sound speed, c_{s2} , and so rapidly catches up with the subsonic compression wave. Using $c_s = \sqrt{\partial p / \partial \rho}$, or for a polytropic gas $c_s = \sqrt{\gamma p / \rho}$, it can be seen that the sound speed, c_{s2} , will be higher in the compressed material. The resulting density ‘spike’ as the two coalesce is known as a *blast wave*.

To further this analysis we consider three phases of a blast wave evolution. To begin with the initial ‘ejected’ mass in the blast wave is greater than the mass swept-up and the blast wave expands freely. Soon the mass swept-up exceeds the ejected mass and the expansion becomes *adiabatic*. Finally as the blast wave decelerates radiative cooling can dominate, and this is known as the *radiative* or *snowplough* phase.

To describe these phases we use a *self-similar* dimensionless solution so that we can compare laboratory and astrophysical shocks in a useful manner via dimensionless parameters such as the Mach and Reynolds numbers. For self-similar blast waves we can model the radius varying as [66]:

$$r(t) \propto \beta t^\alpha \quad (2.33)$$

where β is a constant set by the initial conditions and α is the *deceleration parameter*.

During the adiabatic phase the blast wave evolution is determined by only two variables - the ambient density ρ_1 and the initial deposited energy E . The *Sedov-Taylor* solution for the blast wave radius is used in this phase. Here the swept-up mass is contained in a thin shell behind the shock front and the post-shocked material has negligible mass but high temperature and pressure. The cylindrical solution for β and α can be seen in [21] [67] [68]:

$$r_{st}(t) = f(\gamma) \left(\frac{E}{\rho_1} \right)^{1/4} t^{1/2} \quad (2.34)$$

where $f(\gamma)$ is a gas adiabaticity constant. The cylindrical solution is used here instead of the more common spherical solution as the cylindrical geometry is used in the later experimental chapters.

During the radiative phase the deceleration parameter will change to fit one of two situations. In the *pressure driven snowplough* condition only the blast wave is radiatively cooling and the non-radiating high pressure post-shocked material drives the material with $\alpha = 3/8$. In the *momentum driven snowplough* condition the post-shocked material is also radiatively cooled and the blast wave propagates through it's own inertia, slowing as it gains mass, with $\alpha = 1/3$.

2.3.2 Astrophysical Shocks

Astrophysical systems are typically dominated by strong radiative shocks. A shock is said to be 'radiative' if the *radiative flux* $\geq$ *material energy flux*. For an optically thick plasma the radiation flux approximates to a black body and scales with temperature as $\sim T^4 \sim u_s^8$ where as the material energy flux $\sim T^{3/2} \sim u_s^3$ [69]. Hence strong shocks with a high shock speed, u_s , are more likely to be radiative. Radiation from the shock heated matter can significantly affect the density and temperature upstream. The impact of the radiation will depend on the opacity of the pre and post shocked material. Drake conveniently classes these radiative shocks into four regimes shown in Fig. 2.8 along with astrophysical

examples. A material is optically thick if a photon will have many interactions within

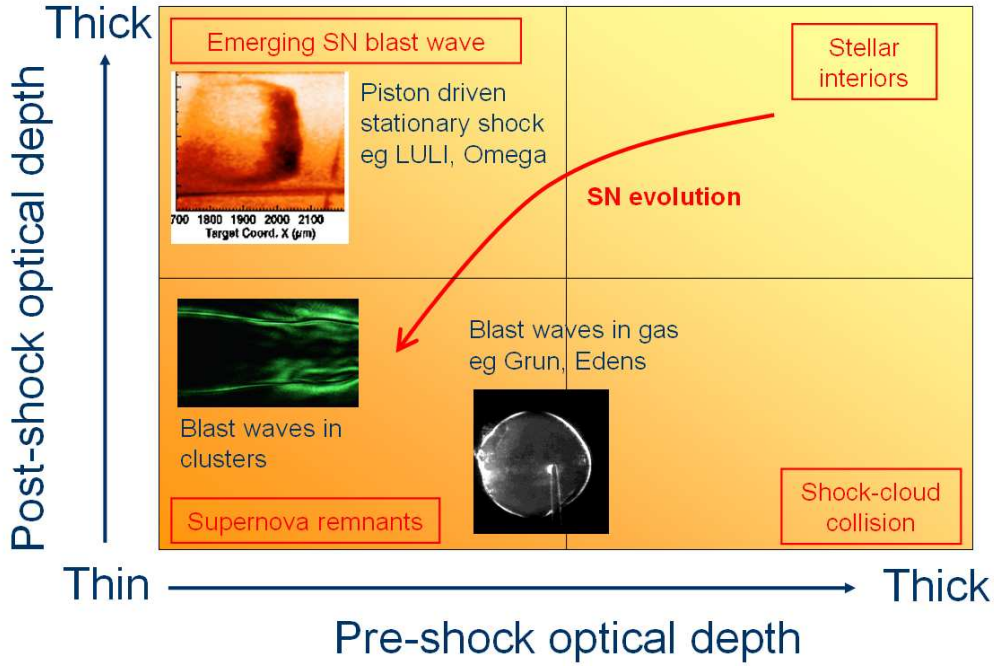


Fig. 2.8: Diagram showing Drake's classification of the four astrophysical plasma radiative regimes by opacity adapted courtesy of Dan Symes. The evolution of an exploding supernova passes through three of these regimes. Some examples are shown of different experiments which access these regimes [70] [71]. Adapted from D. Symes.

it whereas a material is optically thin if a photon will pass through it with very few interactions. A vast majority of observed astrophysical phenomena lie in the 'thin-thin' regime, not least because the escaping radiation makes them easier to detect! Experiments in low density gases access this regime and are discussed in this thesis. Experiments with high Z clusters can also access the 'thin-thick' regime where the pre-shocked material is heated in advance of the shock and ionised by radiation forming a *radiative precursor*.

2.3.3 Requirements for Laboratory Astrophysics Scaling

In order to benchmark numerical models of decade long, light year scale astrophysical scenarios via nanosecond, micron scale laboratory experiments we must ensure we can rigorously scale between the two regimes. The requirements for scaling were first established by Ryutov [1] [5]. The Euler equations for conservation of continuity, force and energy detailed in §2.1.5 are invariant from one system to another under the scale transforms:

$$\mathbf{r}_2 = a\mathbf{r}_1 \quad \rho_2 = b\rho_1 \quad p_2 = cp_1 \quad t_2 = a\sqrt{\frac{b}{c}}t_1 \quad \mathbf{u}_2 = \sqrt{\frac{c}{b}}\mathbf{u}_1 \quad (2.35)$$

where a , b and c are arbitrary positive scalars between systems denoted 1 and 2. Thus the two systems will behave *similarly* and hydrodynamic processes in one system will be identical to the other system if the dimensionless *Euler number*, E_u , is the same [5]:

$$E_u = \mathbf{u}_1 \sqrt{\frac{\rho_1}{p_1}} = \mathbf{u}_2 \sqrt{\frac{\rho_2}{p_2}}. \quad (2.36)$$

In order to carry out the scaling above we have assumed an ideal hydrodynamic fluid. To ensure the validity of the scaling we operate under four key conditions.

Firstly, particles must be localised. Thus the mean free path between collisions must be less than the scale length, L . Ryutov defines the ion-ion mean free path [5]:

$$l_{mfp}^{ii}(cm) \approx 3 \times 10^{13} \frac{T(eV)^2}{n_i(cm^{-3}) \ln \Lambda} \quad (2.37)$$

where $\ln \Lambda$ is the *Coulomb logarithm* which accounts for Debye shielding.

Secondly, heat transport from convection must dominate that from conduction. The *Peclet number*, P_e , is the ratio of convective to conductive heat flow and must be greater than one:

$$P_e = \frac{Lu_s}{\chi_{th}} \quad (2.38)$$

where χ_{th} is the thermal diffusivity. The electron non-magnetic thermal diffusivity is [5]:

$$\chi_{th}(cm^2 s^{-1}) = 2 \times 10^{21} \frac{T(eV)^{5/2}}{n_i(cm^{-3}) \ln \Lambda Z^*(Z^* + 1)} \quad (2.39)$$

where Z^* is the average charge state.

Thirdly, viscosity must be negligible compared to inertial effects. The *Reynolds number*, R_e , is a dimensionless measure of the ratio of inertial to viscous forces and must be greater than one:

$$R_e = \frac{Lu_s}{\nu} \quad (2.40)$$

where the viscosity, ν , is typically dominated by the ion viscosity and for a low Z unmagnetised plasma is [5]:

$$\nu(cm^2 s^{-1}) = 2 \times 10^{19} \frac{T(eV)^{5/2}}{\ln \Lambda A^{1/2} Z^{*4} n_i(cm^{-3})}. \quad (2.41)$$

Finally, the radiation energy flux has to be much less than the hydrodynamic energy flux. Optically thick plasmas will absorb radiation flux and undergo significant changes

whereas thin plasmas will not be affected by the radiation flux. Thus the energy flux due to radiation is highly dependant on the photon mean free path, l_{mfp} . The mean free path can be found as the minimum of the weighted mean free path from Bremsstrahlung radiation in a fully ionised gas, the *Rosseland* mean free path l_{mfp}^{Ross} , or from Compton scattering, l_{mfp}^{Comp} [21]:

$$l_{mfp}^{Ross}(cm) = 1.7 \times 10^{37} \frac{T(eV)^{7/2}}{Z^{*3}n_i(cm^{-3})} \quad (2.42)$$

$$l_{mfp}^{Comp} = 1.5 \times 10^{24} \frac{1}{n_e(cm^{-3})}. \quad (2.43)$$

As described in §2.3.2 in a ‘thick’ plasma the radiative flux is approximately black body. Thus the radiation cooling timescale, τ_{BB} , must be greater than the hydrodynamic time $\tau_{hydro} = L/u_s$ [5]:

$$\tau_{BB} = \frac{\text{energy per unit area}}{\text{black body radiative heat flux}} = \frac{\frac{3}{2}L(n_e + n_i)k_B T}{\sigma T^4} \quad (2.44)$$

$$\tau_{BB}(s) = 1.2 \times 10^{-24} \frac{(Z^* + 1)n_i(cm^{-3})L(cm)}{T^3(eV)}. \quad (2.45)$$

Similarly for a ‘thin’ plasma ($l_{mfp} \gg L$) hotter than a few keV where Bremsstrahlung dominates [5]:

$$\tau_{thin}(s) = 1.4 \times 10^{13} \frac{(Z^* + 1)T(eV)^{1/2}}{Z^* \langle Z^{*2} \rangle n_i(cm^{-3})}. \quad (2.46)$$

As a thin plasma cools line radiation can dominate and τ_{thin} rapidly drops.

Of course this final condition is not desirable for our work as many astrophysical events are radiation dominated and many of the observed phenomena rely on radiation. Radiation causes a precursor to evolve in front of our shocks and, combined with ionisation, lowers the effective adiabatic index γ ultimately causing the shock to decelerate faster. The radiation also results in thinning of the shock front, a higher compression, and enables the observation of instabilities often seen in astrophysics. We can still scale experiments with radiation however if they are carefully designed [72] [9].

2.3.4 Producing an Optically Thin Radiative Blast wave in a Plasma

Previous work on laser produced blast waves in clustered gas show that with a few joule, sub-picosecond light pulse we access temperatures of a few eV over mm scale values at densities of the order of 10^{-4} gcm^{-3} [73]. In Fig. 2.9 we plot this on a temperature against density graph shown by the black dot for partially ionised Xenon gas. Here the solid lines

are based on equations from Drake [74] and Ryutov [5] which use Bremsstrahlung radiative loss in a plasma with local thermodynamic equilibrium (LTE). LTE implies that both the electrons and ions follow the thermal Boltzmann distribution even if the radiation distribution is lower than the blackbody level. Here the green line shows that we fall within the optically thin regime as $\lambda_{Brem} > L$. The blast waves do not fall within the radiatively dominated regime - the red line where $P_{rad} > P_{thermal}$. However, we can still observe radiative effects as the radiative cooling time is less than the cooling time due to convective heat transport, $\tau_{rad} < \tau_{conv}$, again calculated using Bremsstrahlung cooling. Blast waves fall into the radiative flux regime which dominates thermal net flux, $F_{rad} > F_{thermal}$.

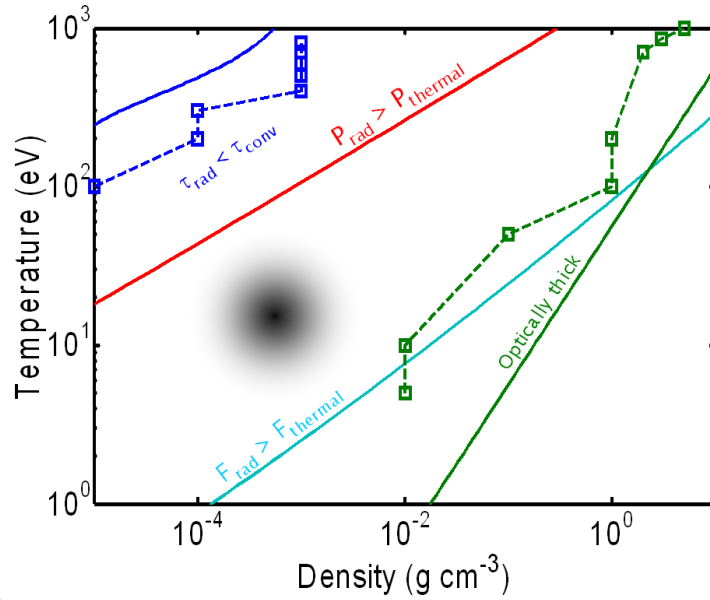


Fig. 2.9: Graph based on Ryutov and Drake (solid lines) showing the boundaries for opacity and the effect of radiation in partially ionised Xenon gas. With new RAPCAL code [75] calculations in collaboration with Symes and Rodriguez which include line radiation and NLTE effects (dashed lines). Cluster blast waves fall approximately within the area shown by the extended black dot.

Symes and Rodriguez have used the RAPCAL code [75] [2] to show that line radiation dominates Bremsstrahlung by more than three orders of magnitude under these conditions. RAPCAL has thus been used to recalculate the regions shown in Fig. 2.9 including effects from line radiation and for a plasma in non-local thermodynamic equilibrium (NLTE), shown as dashed lines. The NLTE calculation accounts for the radiation, electrons and ions in the three different areas around the shock having different temperatures. We clearly see that the cluster blast waves fall in an optically thin radiative flux regime.

2.4 Summary

In this chapter we defined warm dense matter and plasma systems along with some of the methods for modelling them. The inherent difficulties in modelling led us to a discussion of methods and target media for reproducing the necessary temperature and density conditions experimentally. This involved a description of the coupling mechanisms that transfer laser energy into matter in both underdense and overdense regimes. As the laser energy heats the matter we saw that it was vital to understand how electrons and ions are transported, with particular regard to the aspect of a material's stopping power at a given temperature. Knowledge of the ion stopping power in warm dense matter is a key requirement for our experiments and to fully understand particle motion, for instance, in larger planetary cores. Finally we investigated how the natural occurrence of blast waves when high energy densities are rapidly achieved can be scaled to natural astrophysical phenomena using hydrodynamic similarity arguments. This scaling showed that collisionality, heat conduction, radiation and viscosity conditions must be met in the thin-thin plasma regime our blast wave experiments exist within in this thesis.

Chapter 3

Instrumentation and Experimental Methods

This thesis describes the production and diagnosis of the warm dense matter and high energy density plasma states. In order to carry out this work it is necessary to understand several key experimental methods. We begin with a brief description of high intensity laser development which culminates in the Imperial College Nd:Glass and Vulcan chirped pulse amplification lasers used for the majority of work in this thesis. Following the introduction into lasers we look at particle and optical diagnostics used in probing the target media to spectrally, temporally and spatially resolve various properties such as the proton heating beam in warm dense matter experiments and the electron density across our blast waves. We focus particularly on an experimental calibration of proton detectors, radiochromic film and a magnetic spectrometer, using the Surrey Tandetron accelerator. Finally we look at the gas nozzle method for the production of clusters used as the target media in our high energy density plasma work.

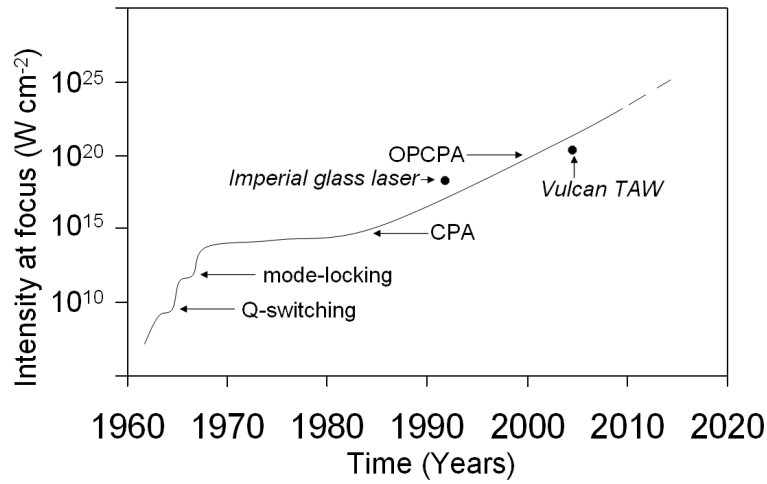


Fig. 3.1: Diagram showing approximate maximum achievable focused laser intensity since the laser's invention in 1960. Jumps in intensity are noticed as the use of Q-switching, mode-locking and CPA overcame various physical limitations.

3.1 High Power Lasers

The first demonstration of light amplification by stimulated emission of radiation (laser) was in 1960 by Maiman [76] with an optically pumped ruby crystal. Since then the challenge has been to produce more powerful lasers that allow access to previously unexplored experimental conditions. Fig. 3.1 shows average increases in laser intensity over time that has relied on four key advances in laser physics: *Q-switch*, *mode-locking*, *chirped pulse amplification* (CPA) and *optical parametric chirped pulse amplification* (OPCPA).

A simple laser scheme contains a gain medium which has been ‘pumped’, typically by an electrical or optical source, to achieve a population inversion. Amplification occurs as a photon with a suitable transition energy stimulates emission from the excited gain medium. Placing a mirror either side of the gain medium creates a cavity which allows further amplification as oscillations occur. The wavelength of the laser is defined by the cavity length, L , such that $\lambda = 2L/n$ where n is a positive integer. Thus the observed wavelengths will be the allowed wavelengths, or modes, that have a suitable transition energy for the chosen gain medium. Using a partially reflective cavity mirror will allow the escape of multi-mode, coherent light.

If we block the cavity until maximum population inversion is achieved by reducing the *Quality factor*, or Q , the full stored energy can be extracted in a single pulse. This *Q-switching* is typically achieved with a Pockels cell and polariser so that light is rejected

from the cavity until the cell fire and rotates the polarisation allowing light to pass through the polariser. Q-switching was introduced in 1961 producing a pulse train of higher energy 120 ns pulses [77].

By preferentially amplifying the modes that are in phase we can lock all the longitudinal modes supported in the cavity in phase so that they oscillate coherently and constructively as a single pulse leading to higher intensities - *mode-locking*. The pulses will be separated by the cavity round trip time and have a duration approximately given by their inverse bandwidth. The bandwidth is $\Delta\nu = N_g c / 2L$ where N_g is the number of modes that experience gain. To achieve this preferential amplification we can use either active or passive methods. Active methods externally and temporally modulate gain in the cavity. More commonly, passive mode-locking relies on the cavity ‘self’ mode-locking via optical effects such as with a saturable absorber mirror (SESAM) which preferentially reflects at higher laser intensities. The Imperial Glass laser uses the SESAM technique.

In the search to reach ever higher laser powers two methods have evolved. One is to create ultra-short pulses, such that durations of ~ 3 femtoseconds are achievable. These can be used in combination with high harmonic generation to achieve 80 attosecond pulses [78] - a decrease of over twelve orders of magnitude in five decades! The shortest pulse limit for propagation is one optical cycle [79]. The other route is to amplify to very high energies through chirped pulse amplification and also the use of multiple beam-lines such as the Petawatt laser at the Lawrence Livermore Laboratories [80]. Here there are two key limitations that prevent us amplifying the pulse *ad infinitum*. Thermal effects caused by energy absorption in optical components can cause damage and reduces the repetition rate of high energy lasers. Typically surface thermal damage occurs if flux levels exceed $\sim 0.16 \text{ J cm}^{-2} \tau^{1/2}$ where τ is the pulse duration in picoseconds [26]. Non-linear optical effects can also cause catastrophic damage. Non-linear optical effects increase as intensity increases so larger beam sizes are used to reduce the risk of damage - increasing the cost.

Non-linear optical effects can best be defined using the Taylor-series description of the polarisation, $\mathbf{P}$, induced by a laser electric field with a constant DC offset, $\mathbf{E} = E_L \cos \omega t + E_{DC}$, in a dielectric material [81]:

$$\mathbf{P}(t) = \epsilon_0(\chi_1 \mathbf{E} + \chi_2 \mathbf{E}^2 + \chi_3 \mathbf{E}^3 + \dots) \quad (3.1)$$

where χ_n is the n^{th} -order non-linear susceptibility. If we look at some of the terms arising from this equation in Table 3.1 they cause some of the common non-linear optical processes

which are relevant to this thesis.

The first term, which we see in all optical systems, is linear with respect to electric field and defines the ‘normal’ materials refractive index, $\eta_0 = \sqrt{1 + \chi_1}$.

The second term causes two well known effects - the *Pockels effect* and *second harmonic generation*. Applying a large DC electric field to a material changes the refractive index via the Pockels effect. This is exploited with birefringent materials in a Pockels cell. Second harmonic generation allows us to convert light to 2ω from ω .

The third term gives rise to the most damaging effect - the *Kerr effect* or *intensity dependent refractive index* such that the material refractive index now becomes $\eta = \eta_0 + I(t) \frac{3\chi_3}{2\eta_0^2 \epsilon_0 c}$. This causes *self-focussing* where a beam with, for example, a higher intensity in the centre will have a higher refractive index at the centre causing the material to act as a positive lens which in turn causes further self-focusing and can lead to catastrophic self-focusing. *Self-phase modulation* also arises from temporal intensity variations leading to temporal refractive index variations which causes a phase-shift which can cause pulse broadening which is useful in creating the bandwidth needed to support ultra-short pulses.

As a useful measure of non-linear optical effects we define the *breakup integral* (or *B-integral*) as a cumulative measure [82]:

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

where L is the optical path length along the z propagation axis. If $B < 1$ non-linear damage is typically avoided. Clearly this limits the length of amplifier material a pulse can pass through and hence the achievable intensity. Chirped pulse amplification aims to overcome this problem and allows further amplification.

Polarisation term	Frequency	Description
$\chi_1 E_L \cos \omega t$	ω	Refractive index
$\chi_2 E_{DC} E_L \cos \omega t$	ω	Pockel's effect
$\chi_2 E_L^2 \cos 2\omega t$	2ω	Second harmonic generation
$\chi_3 E_L^3 \cos \omega t$	ω	Intensity-dependent refractive index

Table 3.1: Non-linear effect only become important for large amplitude electric fields in a dielectric material which are typically only observed with high energy lasers. Listed in this table are some of the important effect which are relevant to this thesis.

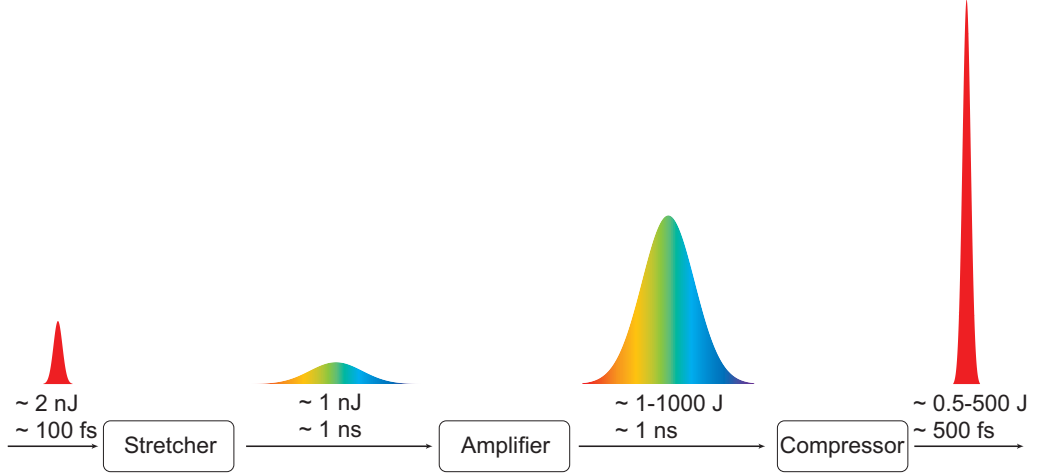


Fig. 3.2: Schematic of the chirped pulse amplification method. A low power ultra-short pulse from the oscillator is stretched in time so that amplification can occur at a lower intensity thus reducing the risk of damage. Compression then shortens the pulse duration to close to the original to achieve high peak power. Gain narrowing often prevents compression to exact original pulse durations which is important in the few femtosecond pulse duration regime. Typical pulse energies and durations are shown for high energy CPA systems.

3.1.1 Chirped pulse amplification

One method of reducing the peak intensity in a laser system which can give rise to B-integral and thermal damage is to increase the beam size and go to larger area and typically more expensive optics. A second method is to increase, or stretch, the pulse duration prior to amplification and re-compress it afterwards. This stretch is carried out by adding frequency dependent phase to the pulse, known as *chirp*. The compression process reverses the chirp. This method using a stretcher, amplifier and compressor as shown in Fig. 3.2 is known as *chirped pulse amplification* and was first proposed in 1985 [83] for optical systems, having previously been exploited in radar as a method for more efficiently extracting energy from an RF amplifier.

The chirp is usually performed with a grating assembly, although other components such as prisms can be used. The added phase as a power expansion around the laser frequency ω_L is [84]:

$$\phi(\omega) = \phi(\omega_L) + (\omega - \omega_L) \frac{d\phi}{d\omega}|_{\omega_L} + \frac{(\omega - \omega_L)^2}{2!} \frac{d^2\phi}{d\omega^2}|_{\omega_L} + \dots \quad (3.3)$$

where $\frac{d\phi}{d\omega}|_{\omega_L}$ is the group delay and $\frac{d^2\phi}{d\omega^2}|_{\omega_L}$ is the group delay dispersion (GDD). Hence gratings disperse frequencies at different angles so they travel a different distance

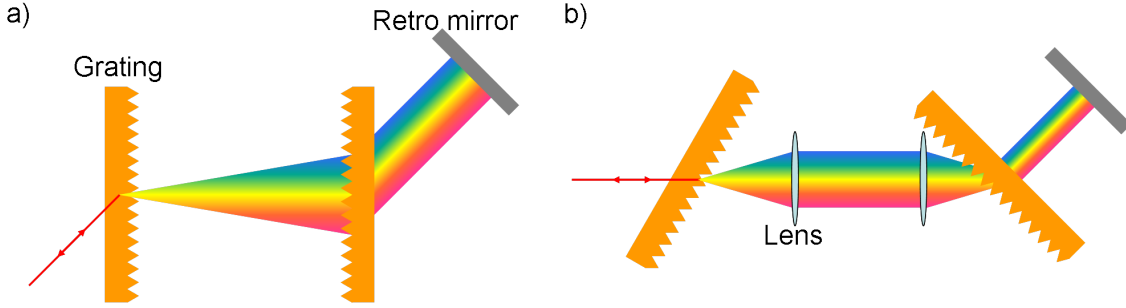


Fig. 3.3: Schematic of (a) a simple parallel grating configuration which adds negative chirp to the pulse and (b) a simple anti-parallel design which adds positive chirp to the pulse. Both configurations are double passed to remove any spatial chirp.

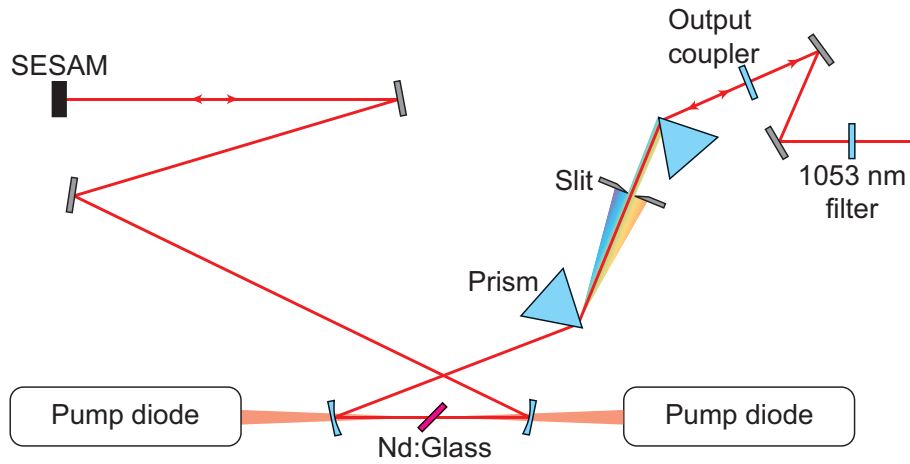


Fig. 3.4: Schematic of the ICLC Glass laser commercial GLX-200 oscillator from Time Bandwidth Products. Two 1200 mW, 800 nm diodes pump the Nd:Glass gain medium which sits in a cavity defined by the partially transmissive output coupler and at the other end by the SESAM which mode-locks the cavity. The two prisms and translation slit are used to tune the bandwidth and central wavelength.

and become spread out in time - a temporal chirp. Fig. 3.3 shows a typical stretcher and grating setup where they are double passed so that any spatial chirp added in one pass is removed in the second, reverse pass. If a positive chirp stretcher is used, where low frequencies lead the pulse, then a complementary negative chirp compressor can be used to reverse the effect. In practice the original short pulse width is not generally fully re-compressed as the amplifier gain medium performs *gain narrowing*. This means that as the gain medium does not uniformly amplify the input bandwidth, the spectrum will be narrowed thus limiting the minimum achievable pulse duration. This effect can be compensated, to an extent, by shaping the input beam spectrum however other higher order phase aberrations can also limit the re-compression.

3.1.2 Imperial College Nd:Glass laser

The Imperial College Laser Consortium (ICLC) neodymium doped glass laser was the first CPA laser built outside America. It was been used by the group for high energy density plasma research since 1994 [85]. It is one of two lasers used for the vast majority of work in this thesis and as such the current configuration, much changed since the original design, is described in this section.

The front-end of the laser is a commercial GLX-200 oscillator from *Time Bandwidth Products* [86]. It produces ~ 200 fs, 1054 nm wavelength pulses at a 100 MHz repetition rate with ~ 2 nJ of energy per pulse. Fig. 3.4 shows the SESAM passive mode-locking design used to produce these pulses. A single pulse is selected using the ‘pre-selection’ and ‘switch in’ Pockels’ cells marked in Fig. 3.5. Before amplification the pulse is stretched to ~ 1 ns duration using a double passed Martinez style reflective anti-parallel grating stretcher. The stretched pulse is then injected into a *regenerative amplifier*. The regenerative amplifier is a linear cavity containing two, flash-lamp pumped, Nd:Glass amplifier heads which are fired once every ten seconds to reduce gain medium cracking from thermal stress effects. The pulse makes multiple passes through the gain medium in the cavity until it is dumped by the ‘switch out’ Pockels cell. A further ‘slicer’ Pockels cell is used to minimise pre-pulse arising from polariser leakage. The regenerative amplifier has a net gain of $\sim 10^6$, reaching an energy of ~ 1 mJ. Spatial filters (SF) and vacuum spatial filters (VSF, used at higher energies to prevent breakdown in air) are used at this point and later in the beam line to improve the beam spatial profile by removing high frequencies in the Fourier plane. At the same time as passing through the spatial filters the beam is also enlarged by a Keplerian telescope to again reduce the B-integral and also to increase the gain volume the beam passes through. The beam is then double passed through an amplifier (~ 10 mJ output) and single passed through two more (~ 100 mJ and up to ~ 3 J). Faraday rotators are used, in conjunction with polarisers, after the second and third amplifiers to optically isolate the system and prevent any back reflections being re-amplified and causing damage. Leaving the final amplifier, which fires once a minute, the 50 mm beam is re-compressed in a double-passed Treacy style parallel 1740 lines/mm grating compressor to ~ 700 femtoseconds with ~ 50 % throughput.

A cross-calibrated pyroelectric energy sensor is used at the zero-order reflection of the first compressor grating which contains 16% of the energy of the compressed beam. Running the regenerative amplifier in an unsaturated regime results in significant shot-

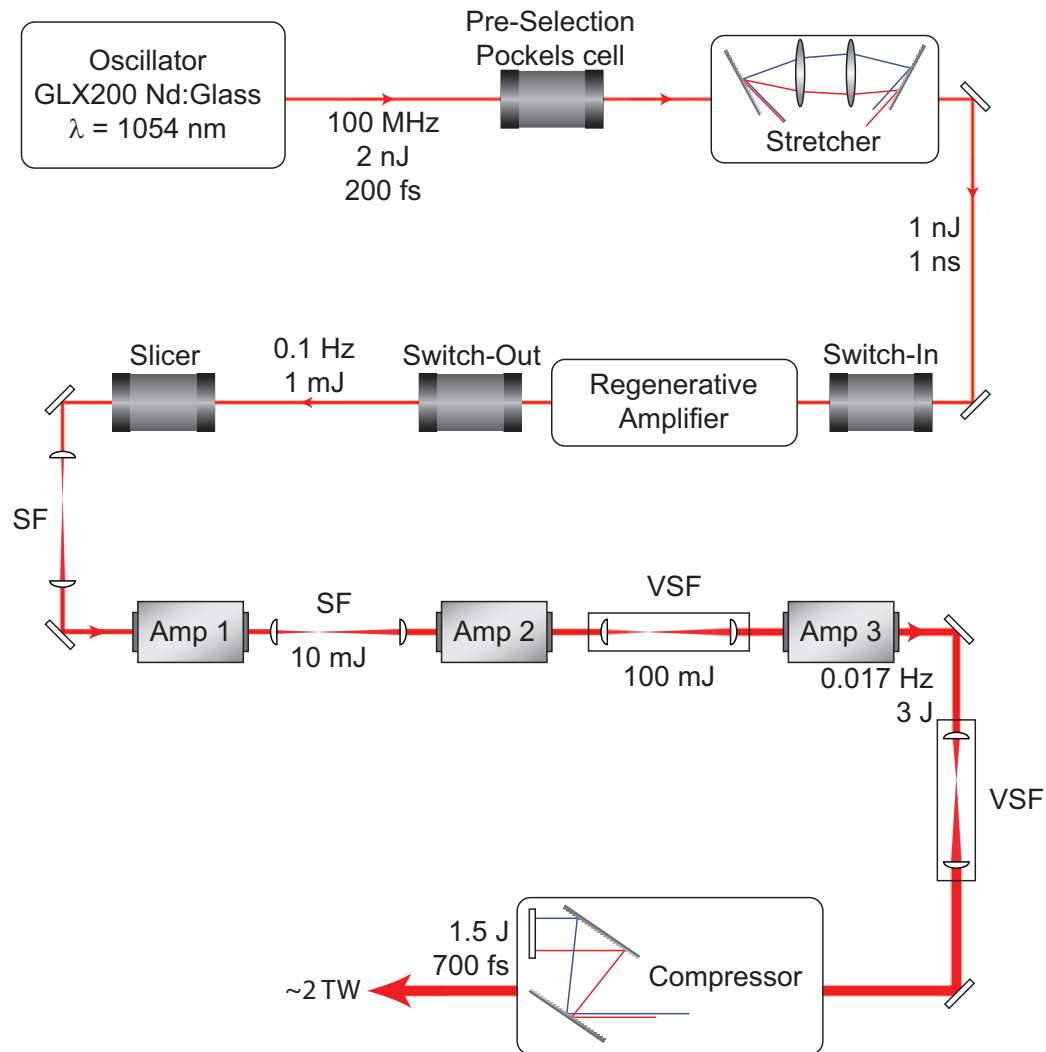


Fig. 3.5: Schematic of ICLC Nd:Glass laser. The pulse train from the oscillator is stretched in a double passed stretcher and amplified after one pulse has been selected by the pre-selection and switch in Pockels' cells. The amplified pulse is re-compressed via a double passed, opposite chirp compressor grating assembly.

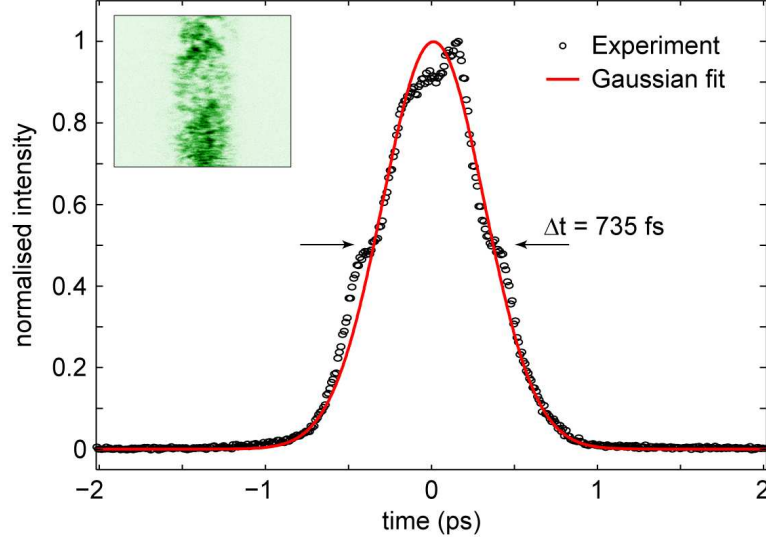


Fig. 3.6: Inset shows raw data of 527 nm signal from the single shot autocorrelation. Data is integrated vertically to achieve the black data points. A Gaussian fit provides a pulse duration measurement of 735 femtoseconds. Adapted from [28].

to-shot energy instability - a target average energy of 1 J has a FWHM energy spread of $\sim 270 \text{ mJ}$. While this limits repeatability our relatively high repetition rate enables us to scan over a large range of energies during an experiment.

Pulse duration is measured on a single shot basis using a second-order autocorrelator [87]. A beam sampler takes the original pulse which is split in two and then overlapped temporally and spatially at an angle in a second-harmonic crystal. Second harmonic generation occurs where the two pulse replicas overlap and the temporal structure of the original pulse is mapped onto the spatial structure of the green 2ω pulse as seen on a camera. Thus the pulse duration can be recovered (making an assumption of a Gaussian pulse shape) as shown in Fig. 3.6.

Measuring the focal spot size at full power is non-trivial. Imaging the focal spot with an alignment or low power beam as shown in Fig. 3.7 is a good first order measurement. However, for a full energy shot wavefront distortions may occur which affect the focusability on a shot to shot basis. In foil interaction experiments laser pre-pulse may also affect the target position, further changing the focal spot. Ideally *in situ* measurements of focal spot size on every shot should be carried out. Imaging can not be carried out at the fundamental wavelength because scattering and plasma emission enlarge the measured spot size. It has been shown that 3ω and higher harmonic imaging are an accurate and reliable measurement of *in situ* spot size [88]. These measurements are particularly important for interactions where a small, clean focal spot is vital to achieve high intensities and may

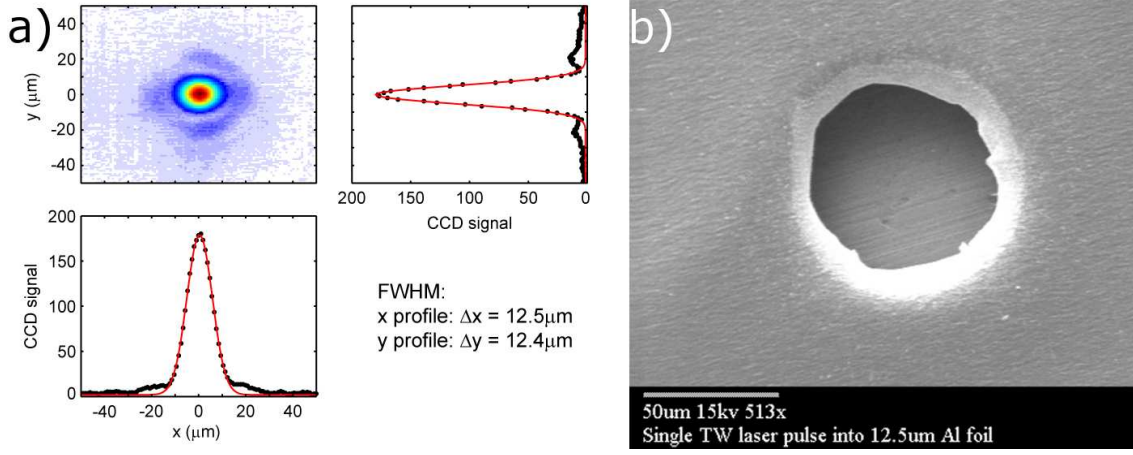


Fig. 3.7: a) Vertical and horizontal lineouts from a magnified image of the low power CW focal spot show a best case focal spot size of $\omega_0 = 12.5\mu\text{m}$. The $M^2 \approx 1.1$ measures how near to the diffraction limited case ($M^2 = 1$) the laser is [89]. Figure courtesy of Dr Matthias Hohenberger. b) Scanning electron microscope image of the hole remaining after the Glass laser has hit a foil for proton production at near to best focus. The $50\mu\text{m}$ hole is larger than the CW focal spot and provides an upper limit for the focal spot size on a full shot. While most of the size difference can be accounted for by long time scale thermal damage after the laser pulse has arrived it is a clear demonstration of the difficulty in definition of focal spot size.

also be used as a qualitative measurement of pre-pulse.

Using the measurements described above we can gain an order of magnitude estimate of the peak Gaussian laser intensity at focus with [84]:

$$I_0 = \frac{E}{\tau\pi\omega_0^2} \sqrt{\frac{4\ln 2}{\pi}}. \quad (3.4)$$

Thus we access peak intensities $\leq 10^{18} \text{Wcm}^{-2}$.

3.1.3 Optical Parametric Chirped Pulse Amplification

OPCPA operates on the same underlying concept as CPA except the gain method used is an *optical parametric amplifier* (OPA). It was first experimentally demonstrated by Dubietis in 1992 [90]. OPA relies on *difference frequency mixing* where a *pump* photon of frequency ω_1 decays in a dielectric medium to two photons with an arbitrary combination of frequencies ω_2 and ω_3 such that $\omega_1 = \omega_2 + \omega_3$ as shown in Fig. 3.8 [81]. To define the temporal and spectral structure of the resulting combination of frequencies a small *signal* input of ω_2 results in an output amplification of ω_2 . The ω_3 pulse is known as the *idler* and can be useful in its own right, though generally represents an energy loss mechanism.

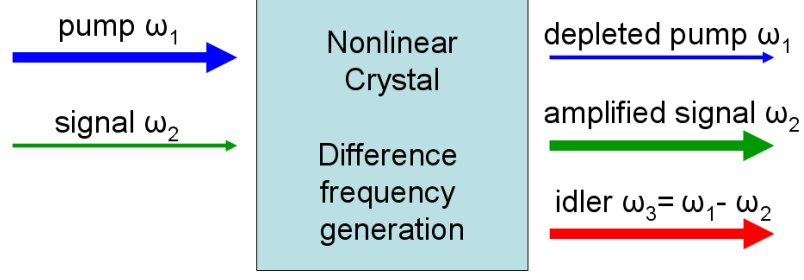


Fig. 3.8: Schematic showing the mechanism for optical parametric amplification. The pump excites the medium to a virtual state which then decays via another virtual state to ground. This two stage decay results in the signal and idler photons.

Theoretical efficiencies of pump to signal conversion are $\sim 50\%$ as a significant fraction of the pump is converted to idler. Higher repetition rates than conventional amplifiers are achievable as there is very little heating of the non-linear medium since the pump energy is fully converted to the signal and idler, and no energy is stored in the system.

OPA's are not limited to the conventional gain bandwidth of a gain storage amplifier and as such do not suffer from gain narrowing in the same way. The bandwidth of an OPA is determined by the length and dispersion of the non-linear crystal. For efficient amplification phase matching must be tuned to the signal wavelength. The phase mismatch, Δk can be defined by:

$$\Delta k = k_1 - k_2 - k_3 \quad (3.5)$$

where k_i are the respective wave-vectors.

Since amplification only occurs during pumping no energy is stored in the gain medium which means that a very high contrast and intensity can potentially be achieved. However, an equivalent of amplified spontaneous emission, so called *parametric fluorescence* can create pre or post pulses.

3.1.4 The Vulcan laser in target area west at the Rutherford Appleton Laboratory

The Vulcan laser at the Central Laser Facility of the Rutherford Appleton Laboratory was used for high energy scaling of our blast wave work and for the majority of the warm dense matter work described in this thesis.

The Nd:glass laser has eight long pulse ~ 2.6 kJ, ~ 2 ns beams and one short pulse ~ 700 J, ~ 500 femtosecond beam. In 2005 the Vulcan laser held the record for the highest

focused intensity on target in the world of 10^{21} Wcm^{-2} .

The laser has two main target areas with different capabilities, Target Area Petawatt (TAP) and the area where most of our work was carried out Target Area West (TAW). TAW received an upgrade between two of the experiments in this thesis allowing higher achievable intensities with the use of a dielectric grating. Fig. 3.9 shows the beam-line for these two target areas which rely on a combination of CPA and OPCPA technologies described earlier.

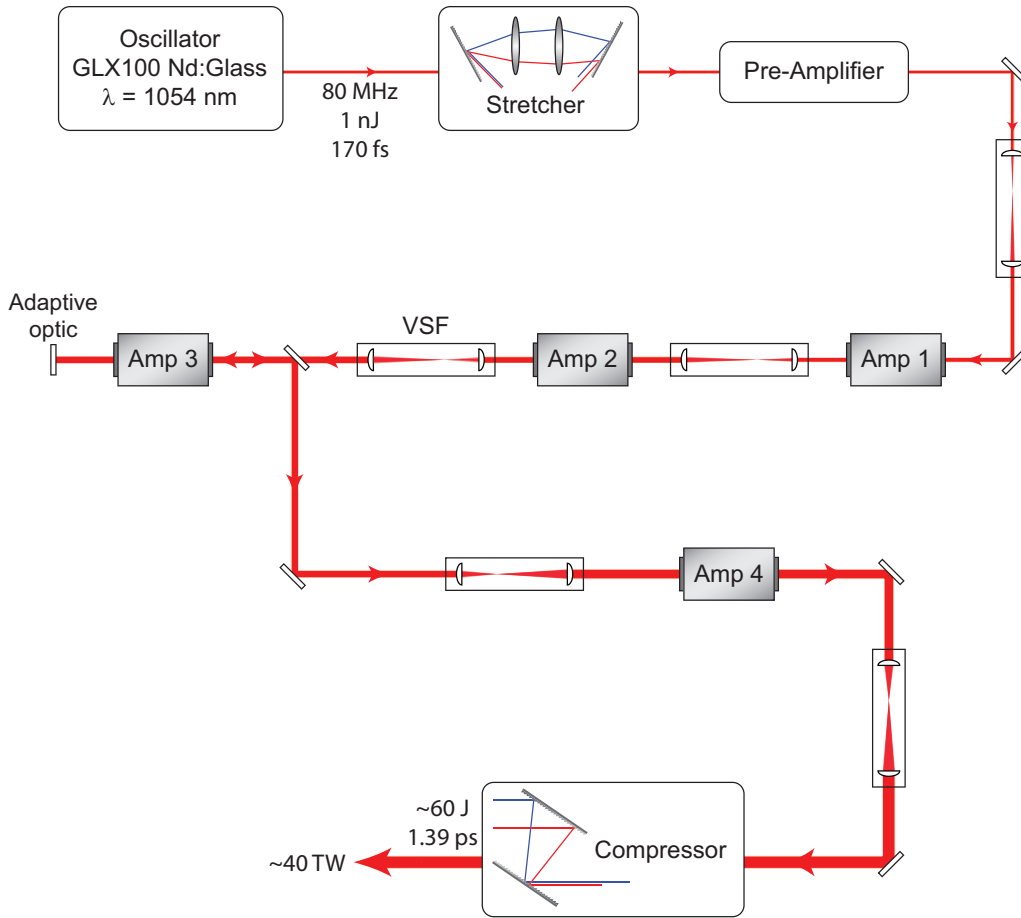


Fig. 3.9: Schematic of Vulcan Target Area West beam-line pre 2009 upgrade. Pulses from the oscillator are stretched so that amplification can be carried out at lower intensities to reduce the B-integral and risk of damage. After several amplification stages and intermediate vacuum spatial filters the beam is re-compressed to ~ 1 picosecond. An adaptive optic, or deformable mirror, is used to correct for some spatial phase aberrations to improve the focusability.

3.2 Ion Diagnostics

In order to diagnose an ion beam four key properties have to be characterised. These are the number, species, position and velocity distribution of the ions. The velocity vector will enable us to recover the ion energy spectrum and the beam divergence. In most practical experiments producing proton beams there will also be higher Z ions within the beam and so it necessary to diagnose the charge and mass of the captured particles.

Ion detection techniques can be split into two categories - materials that will detect ions and the tools they can be used in conjunction with to further characterise the ions. In both of these categories we need to consider the requirement for resolution of energy, species, dynamic range and spatial profile.

3.2.1 Materials that detect ions

Two types of solid state ion detector are used within this thesis - CR-39 and radiochromic film (RCF). Image plate (photostimulable phosphor plate) is a third type of solid state detector but requires special scanning equipment to use and as such does not feature in this thesis. Scintillators in combinations with cameras or micro-channel plates (MCP) are a growing technology for high energy ion detection as they allow high repetition rate readout however they are not used in this thesis as they are largely un-calibrated at low (few MeV) proton energy, prone to noise and in the case of MCPs require very low vacuum ($< 10^{-6}$ millibar).

CR-39 is a clear, colourless thermosetting plastic commonly used in sunglasses and lenses. It has a useful property in that it can record nuclear particle tracks but is mostly insensitive to electron and photons. It was discovered in 1978 at the University of California by Cartwright, Shirk and Price. They found CR-39 will record particles with $Z/\beta > 6$ [91] where β is the relativistic velocity meaning it will record protons up to 13 MeV. It has the chemical formula $C_{10}H_{18}O_7$ and a density of $\sim 1.30 \text{ gcm}^{-3}$. A charged particle entering the plastic will be brought to rest by losing energy through ionisation, damaging the polymer. An ion needs to have a minimum of $\sim 75 \text{ keV}$ to cause this damage. If the CR-39 is etched in a $\sim 80^\circ\text{C}$ bath of 6M sodium hydroxide the area where the chemical bonds are damaged will be preferentially etched. The final size of these etched tracks will depend upon the depth at which the ion is stopped (its initial energy), the ion species size and the etching time. The smaller the pit the higher the energy ion. These $\sim$ micron size pits can be observed with a standard white light microscope as seen in Fig.

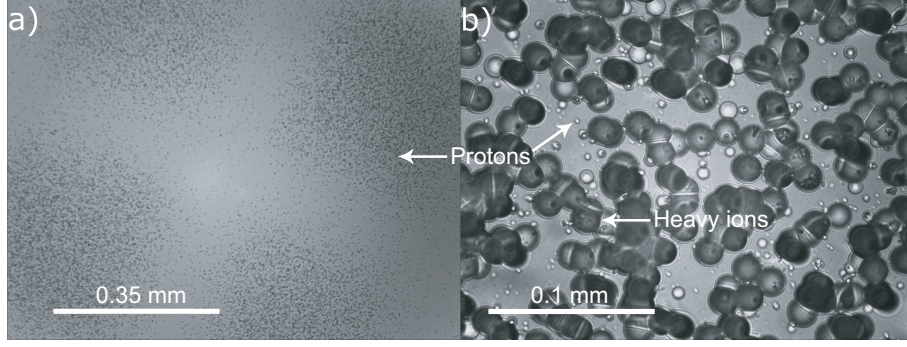


Fig. 3.10: a) Microscope capture of unsaturated CR-39 tracks, each track corresponds to an individual ion hit. The protons result from a laser foil interaction on the Imperial Glass laser system. The resulting imprint from a mesh can be seen which shows different levels of CR-39 signal, as the tracks overlap they can be difficult to distinguish. The plastic was etched in NaOH solution for one hour. A standard white light microscope at $\times 40$ magnification was used. b) CR-39 tracks etched for 4 hours from a Vulcan target area petawatt experiment at $\times 100$ magnification. Different size tracks correspond to different species of ions - typically small tracks are protons and the larger tracks are heavier ions.

3.10. A key advantage of CR-39 is that the primary measurement, the number of tracks, is sensitive to fluence not to dose - simplifying analysis. However, counting tracks can be time consuming, even using automated counting techniques, and ultimately limits the dynamic range as saturation occurs when individual tracks cannot be resolved, typically at fluences of $\sim 10^6 \text{ cm}^{-2}$.

Radiochromic film (RCF) is often used in conjunction with CR-39. RCF is a thin plastic film that responds to all ionising radiation but only weakly to optical light. An aluminium foil filter is typically used on the front of the RCF to block heavy ions and optical radiation. Fig. 3.11 shows the structure of a typical make of RCF, HD-810, used throughout this thesis. On exposure to ionising radiation the dye changes colour from transparent to blue by a photo-chemical reaction. The resulting optical density is proportional to the dose of radiation incident. Dose, measured in grays, is the energy absorbed per unit mass. We can describe this in more useful observables:

$$D = \frac{\text{Joules}}{\text{kilogram}} = \frac{\varphi(T - T_0)}{\rho r} \quad (3.6)$$

where φ is the proton fluence, T is the energy on exiting the active region, T_0 is the energy on entering the active region, ρ is the film density and r is the range in the active layer.

For accurate absolute measurements of proton dose calibration can be carried out on an accelerator so that the optical density measured using a standard scanner can be correlated to a well-known proton dose - this is vital at proton energies $< 1 \text{ MeV}$ as the protons do

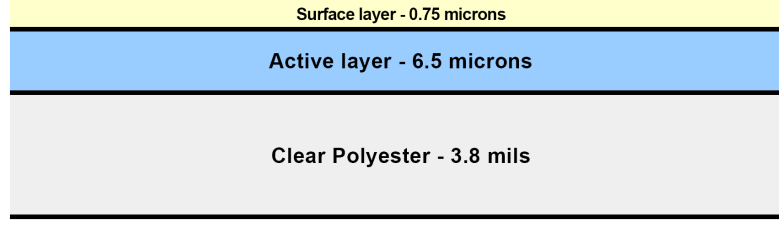


Fig. 3.11: Schematic of HD-810 radiochromic film from Gafchromic® [92]. The active layer is triphenyl methane dye cyanide.

not penetrate the RCF fully and there is a lack of literature showing calibration at these energies. Once an accurate dose is known a fluence can be calculated for a given proton energy. For 1 MeV protons HD-810 can detect fluences between approximately $1.3 \times 10^8 \text{ cm}^{-2}$ and $1.3 \times 10^{11} \text{ cm}^{-2}$. The dynamic range is to a large degree dictated by that of the scanner used.

We calibrated HD-810 RCF using the 2 MV Tandetron accelerator from High Voltage Engineering Europe at the University of Surrey Ion Beam Centre. The Tandetron generates proton beams up to 4 MeV which are highly mono-energetic ($\pm 10 \text{ keV}$). Using a picoammeter to measure current we exposed the RCF to different known fluences and energies to obtain the results shown in Fig. 3.12 and Fig. 3.13. The RCF was scanned using a standard flatbed scanner with film attachment (CanoScan LiDE 600F) and a high dynamic range Artixscan 1800F transmission scanner. A Stouffer calibrated optical density chart was also scanned in order to convert pixel value to optical density. Whereas much of the literature uses red, green and blue pixel values and different doses we use an average of the values in order to simplify analysis with very little loss of dynamic range.

Fig. 3.12 shows that optical density varies with energy for a constant fluence ($\sim 5.6 \times 10^9 \text{ cm}^{-2}$) in two approximately linear regimes. It increases linearly with energy where the protons' range is within the active region but as expected response peaks when the protons' Bragg peak is beyond this region and then decreases linearly with energy beyond $\sim 1000 \text{ keV}$. The calculated dose suggests the optical density should continue to increase in the active region towards the origin. The author attributes this discrepancy to the protons' reduced range in the active region resulting in less of the dye being exposed. In other words if the proton only travels through half of the active region only half of the dye will change colour resulting in a lower optical density change than expected. The proton range is calculated using the stopping range of ions in matter code SRIM [57]. RCF is often used for dose measurement during medical therapies. The medical literature for RCF

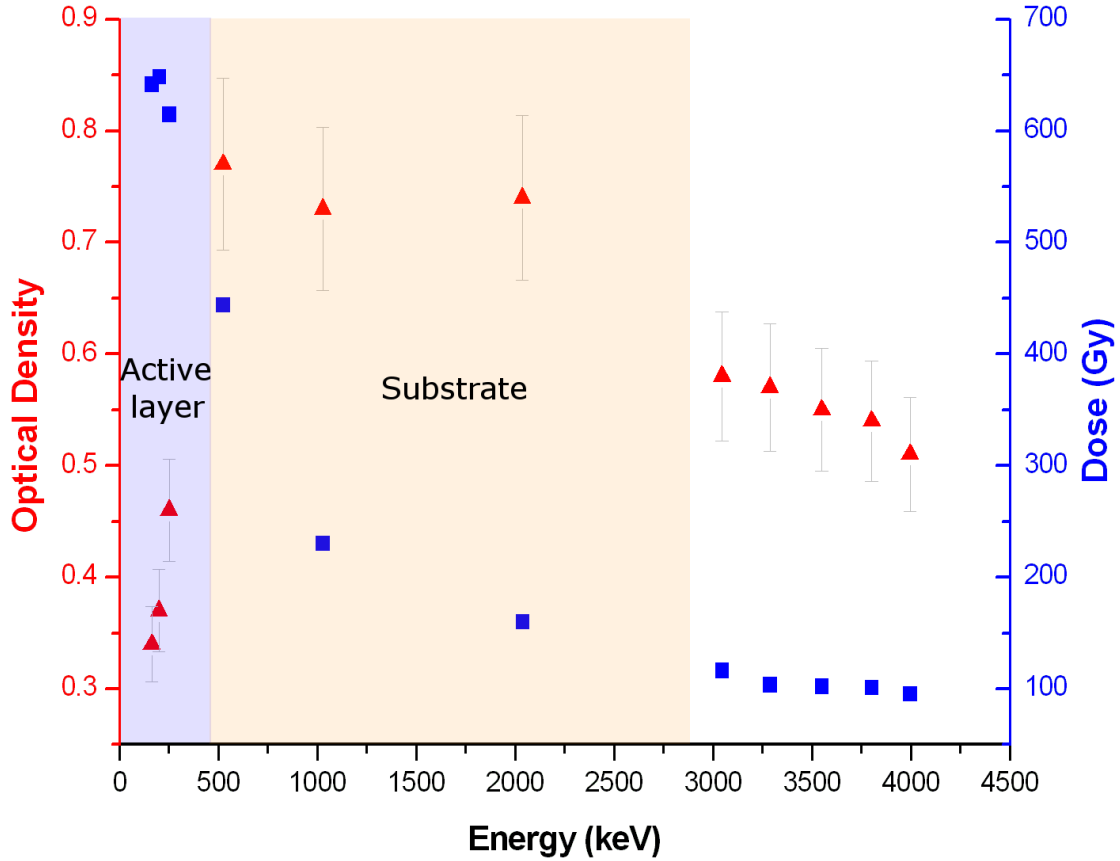


Fig. 3.12: Results from the Surrey Tandetron for RCF calibration of proton energy against measured optical density using a calibrated (with a Stouffer step filter) standard flatbed scanner. The error in energy is ± 10 keV. The fluence is held constant at $\sim 5.6 \times 10^9$ cm⁻². Calculated dose is shown by the square blue points. The shaded regions show the active layer and substrate of one piece of RCF the protons will penetrate through to as calculated using TRIM. The data falls into two approximately linear regimes - inside and outside the active region. In the active region optical density decreases even though the dose increases as a result of all the proton energy being deposited in the dye. The author attributes this difference to only a portion of the dye being exposed to the protons.

suggests that the optical density change is largely independent of energy. My data shows that this is only true where the proton range is much greater than the RCF thickness. Where RCF is used in a stack this condition does not necessarily hold true. Although the literature shows several similar RCF calibrations [50] [93] [94] [95], the author believes this is the first time the energy calibration has been carried out at sub-MeV proton energies and thus the first observation of this effect.

If we hold the energy constant at 1.03 MeV and increase the fluence we observe a linear increase in optical density as seen in Fig. 3.13 within the limits dictated by the dynamic range of the scanners. Here the dose is directly proportional to fluence.

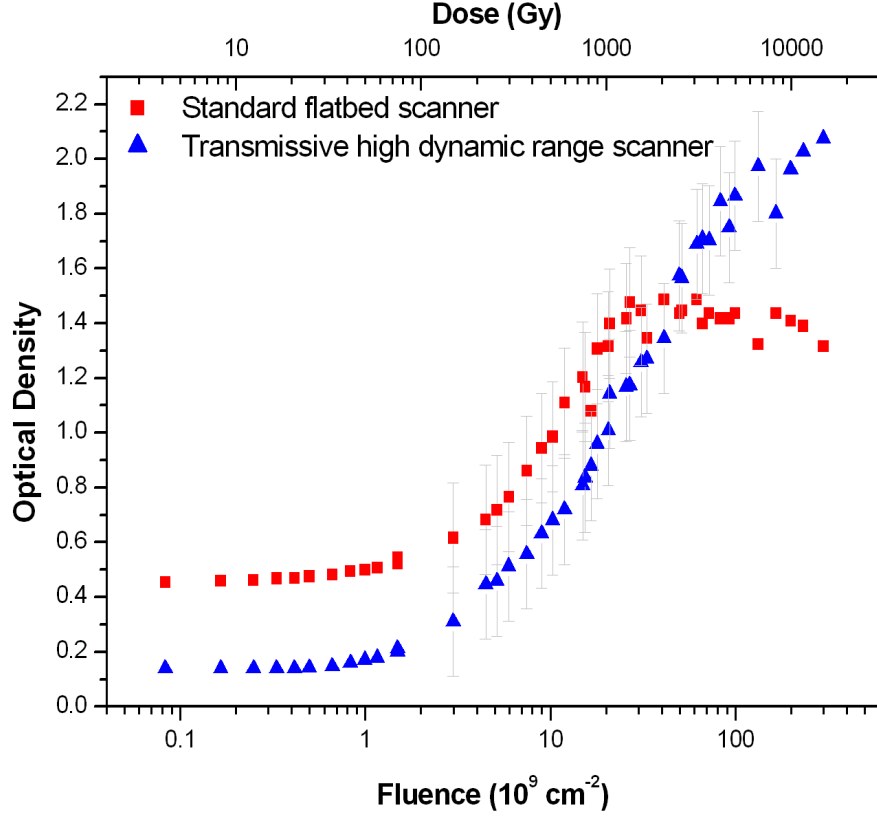


Fig. 3.13: Results from the Surrey Tandetron for RCF calibration of proton fluence against measured optical density using a calibrated (with a Stouffer step filter) standard flatbed scanner (red squares) and a high performance transmission scanner (blue triangles). Error bars are only shown in the linear region as above this the scanner saturates and below this again the dynamic range of the scanner limits the detection ability. The energy is held constant at 1.03 MeV meaning that the dose varies linearly with fluence.

Using the linear relationship of optical density to energy beyond 1000 keV for constant fluence and the linear relationship to fluence for constant energy within the limits dictated by the scanners dynamic range we can solve the simultaneous equations to give a relationship between optical density and fluence for a known energy:

$$\varphi = OD + (T7.769 \times 10^{-5}) + 0.53127 \quad (3.7)$$

where OD is optical density and T is the initial proton energy entering the film. Here the assumption is that the linear fluence relationship holds for any given energy. This equation is for the standard flatbed scanner. The uncertainty in fluence for a given optical density is $\pm 15\%$ based on the experimental data and is dominated by the accuracy, or lack thereof, in the calibration fluence measurements.

3.2.2 Detector stacks

A powerful use of the RCF and CR-39 detectors is to combine them into an imaging detector stack. The materials' proton stopping power means that for every layer of the stack an energy bin can be assigned to the protons that stop in it and hence provide the majority of the dose in that layer. Fig. 3.14 shows an example stack with the corresponding energy bins calculated using the Monte Carlo code TRIM [57]. A stack can be used to calculate beam divergence and also carries some temporal information. Inherently the higher energy protons which penetrate deeper into the stack will have travelled through the interaction region faster, and so will carry information from an earlier time than the less penetrating protons.

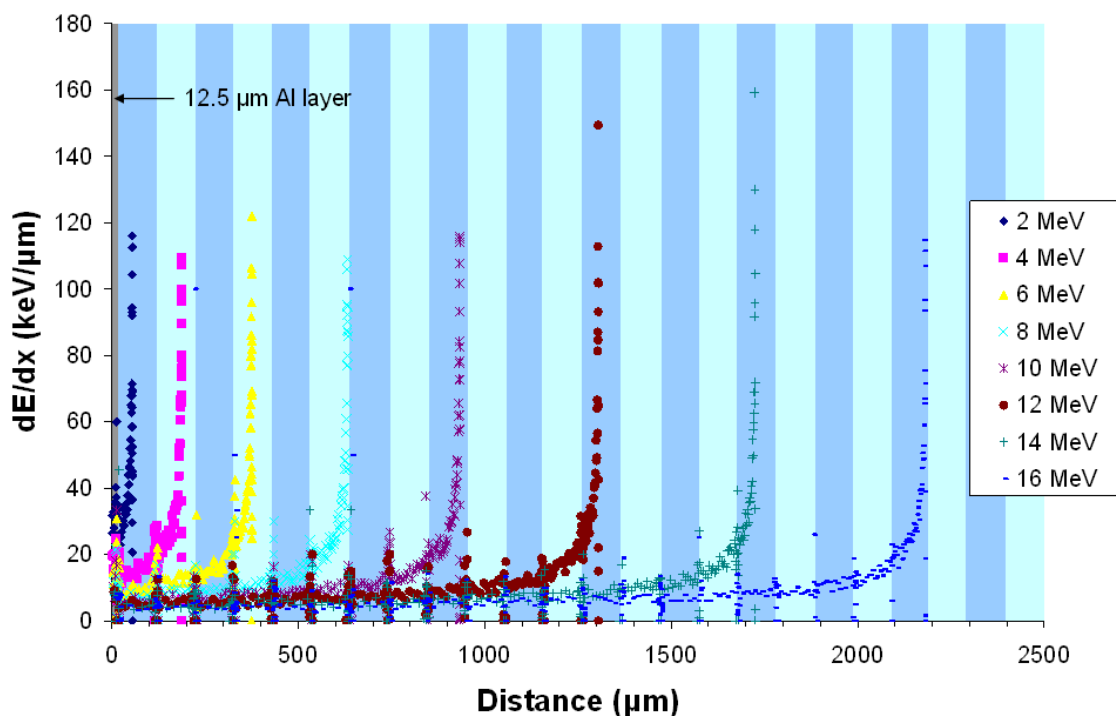


Fig. 3.14: Typical RCF stack with proton energy loss calculated by the author using TRIM for several proton energies. A front layer of aluminium filters visible light from the RCF. In this stack 24 layers of RCF are used, other variations involve using CR-39 or metal filters for some of the layers. The small peaks at the layer boundaries are an artefact of the code which arises from a mono-layer size binning error at the layer edges.

Measuring the energy spectrum from a stack is not very accurate. The large energy bins mean that calculating fluence from the observed dose requires an average energy to be used. Also, although the Bragg peak of a proton occupies a relatively short distance at the end of the range, there is still energy deposition throughout the particles trajectory.

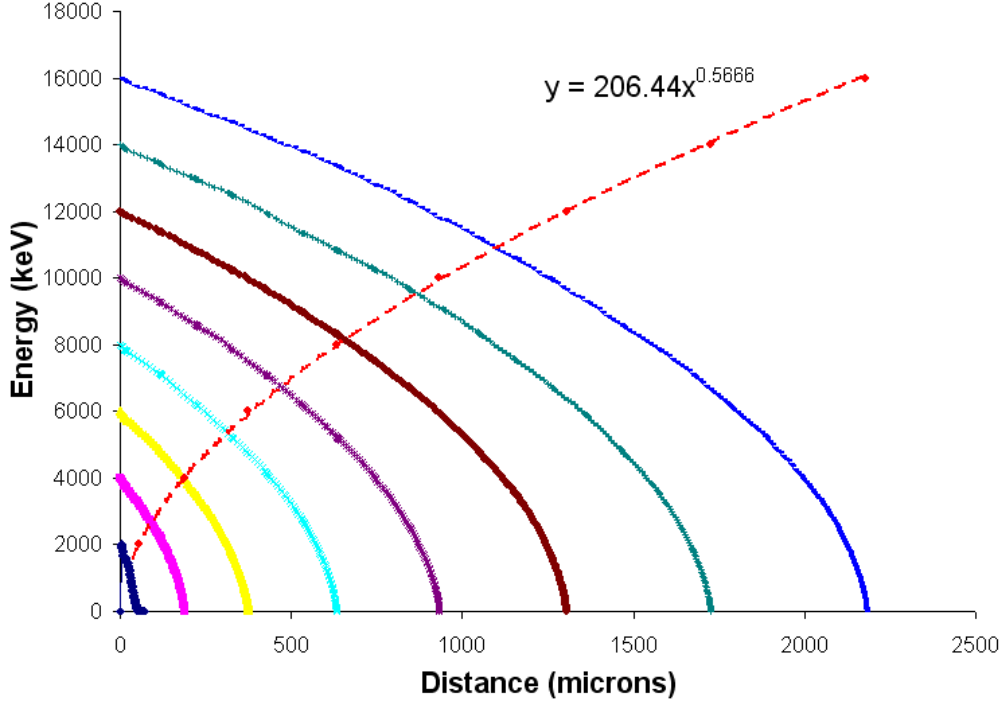


Fig. 3.15: Typical RCF stack with several proton energy ranges calculated by the author using TRIM. Curve fitting software was used to find an equation for the proton range through the stack shown as the dashed line of initial energy against range. It should be noted that this range is an average as stopping is a statistical process involving multiple collisions.

Thus the dose from protons that stop in later energy bins must be subtracted to obtain an accurate fluence. This deconvolution can be carried out by working from the back piece of the stack forward as the dose at the back will not have any contribution from lower energy protons that stop elsewhere in the stack. Fig. 3.16 shows data taken during the RCF calibration on the Surrey Tandetron where a small stack was used to demonstrate that protons stopping in a later piece of RCF deposit a significant dose in earlier pieces. The dose was obtained from the fluence measurement derived from optical density using Eq. (3.7).

3.2.3 Low energy magnetic spectrometer calibration

To measure the ion energy spectrum more accurately, a magnetic spectrometer was built to detect ‘low’ energy protons used for warm dense matter experiments. The concept behind the magnetic spectrometer is simple. Two large permanent magnets ($\sim 0.8\text{ T}$) establish a magnetic field in the z -direction. An ion travelling in the x -direction will pass through the entrance slit, which to a large extent defines energy resolution, to be dispersed in the

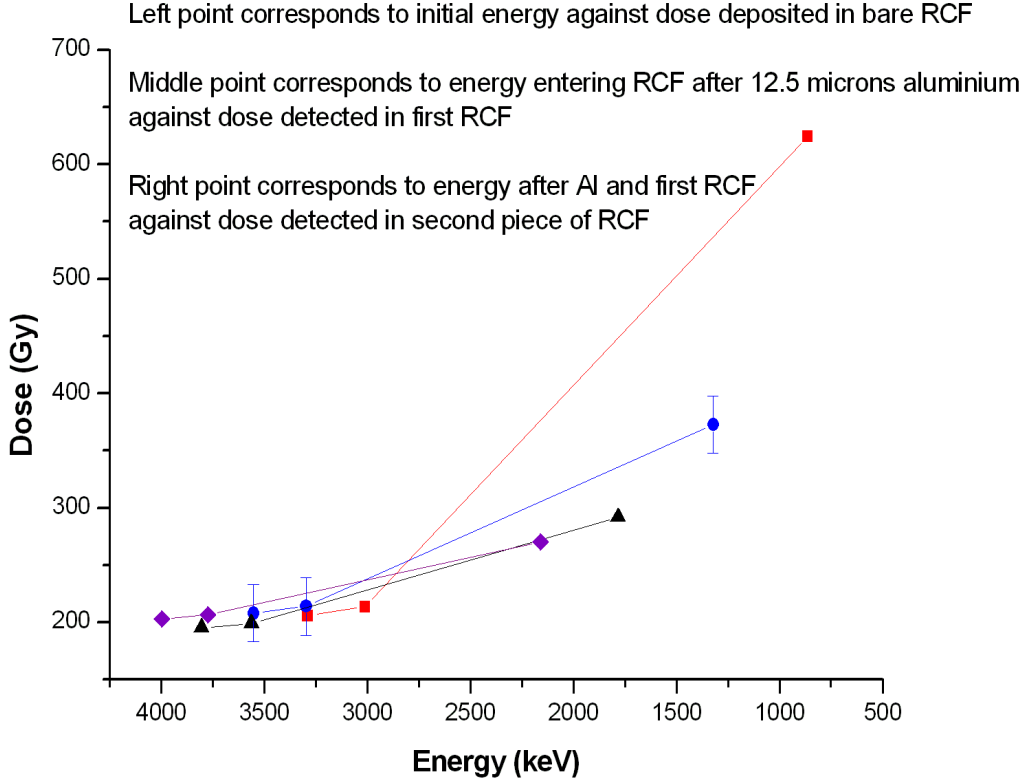


Fig. 3.16: Using the Tandetron accelerator at Surrey a small RCF stack consisting of a 12.5 μm aluminium mask and two layers of RCF was calibrated against energy for approximately constant fluence ($5 \pm 1 \times 10^9 \text{ cm}^{-2}$). A limited range of energies were used due to time and equipment constraints. Four initial proton energies were used: 3.29, 3.55, 3.80 and 4.00 MeV (diamond, triangle, circle and square respectively). For each energy three values were obtained: the dose on a single piece of RCF, the dose on RCF after the aluminium mask and finally the dose on the second layer of RCF having passed through the mask and the first layer (left to right respectively). The data shows that within a stack a proton will still deposit a significant dose within the layers before it's termination point. However, the dose in the preceding layers is only weakly dependent upon energy as can be seen by the approximately constant dose for the first two data points in all three data sets.

y -direction by the Lorentz force, $\mathbf{F} = q\mathbf{v} \times \mathbf{B}$. The fastest, heaviest and least charged ions will be deflected least. Intuitively, the fastest ions might be expected to be deflected most as the force is greatest but this is countered by the fact that they spend the least time in the field. For a constant $\mathbf{B}$ field the ion trajectory is along the gyroradius, $r_g = mv_{\text{perp}}/qB$, where v_{perp} is the velocity component in the x -direction. However, 'fringe' effects result in a non-uniform magnetic field and so calibration is required for accurate results. Initial $\mathbf{B}$ field calibration was carried out with a Hall probe and a particle tracking model. However for these low proton energies small inaccuracies in measurement of the field or geometry can result in large shifts in calculated energy. Calibration was also carried out using Americium 241 as a known alpha source at 5.5 MeV. This was not very accurate however as the internal

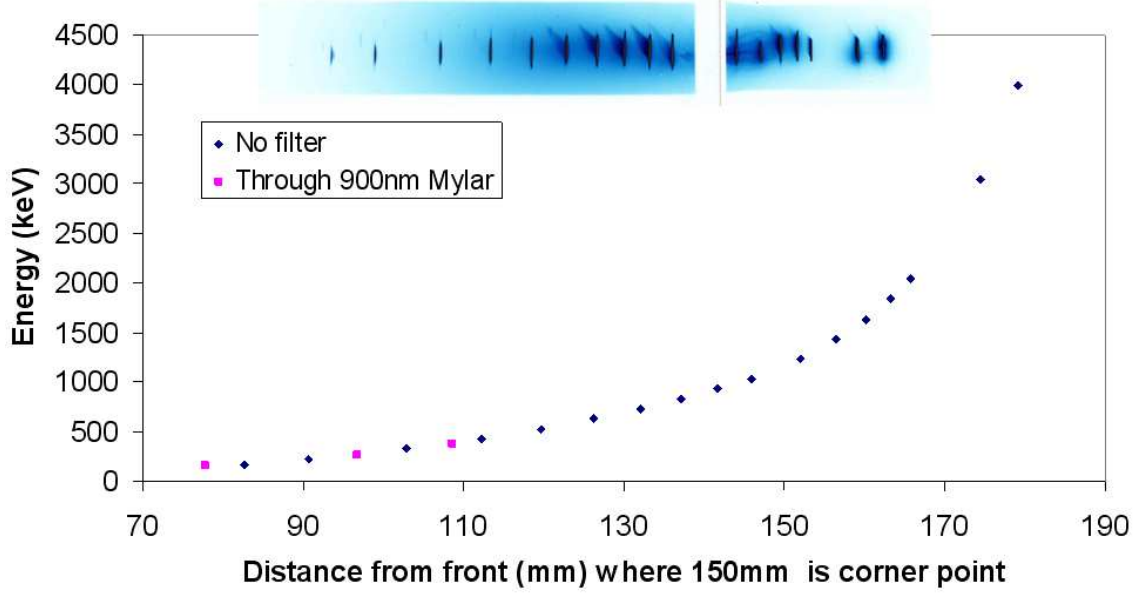


Fig. 3.17: Results of magnetic spectrometer calibration. This figure allows accurate mapping of a point on the detector plane to a specific proton energy. The error in both distance and energy is $< 1\%$ and is covered by the data points. The inset shows the raw data. The accelerator could not achieve energies below 80 keV so a Mylar film with a known stopping power was used to achieve lower energies (pink squares). Lower energies could not be achieved as the energy spread was too wide after traversing the Mylar film for accurate measurement.

stopping in the source resulted in a broadband alpha emission. Yet another method for calibration uses a laser generated proton source with a known filter on the detector plane. Using SRIM we know that a given thickness of aluminium will only transmit protons above a certain energy. However, this results in further inaccuracies as there is a threshold for detection on the detector material. Because none of these calibration methods agreed to a reasonable accuracy we took the spectrometer to the Surrey Tandetron where a known mono-energetic proton beam was used. The results for this are shown in Fig. 3.17 which allows us to fit a calibration curve, $f(x)$. In analysing spectra recorded in the spectrometer we measure fluence, φ_{meas} , per length, Δx , on the detector plane. Clearly a given unit length represents a different unit of energy, ΔE , depending on the dispersion curve. To correct for this the correct fluence, φ_{corr} , is:

$$\varphi_{corr} = \frac{\varphi_{meas}}{\frac{df}{dx} \Delta x}. \quad (3.8)$$

To differentiate between protons and heavier ions aluminium filters on the detector plane were used which transmitted protons but stopped higher Z ions.

A quantitative derivation of the magnetic spectrometer error can be found in the appendix.

3.2.4 Thomson spectrometer

A Thomson ion spectrometer, also known as a Thomson parabola, is similar in operation to the magnetic spectrometer but also adds an electric field and uses a pinhole instead of slits. The electric field is supplied by setting up a potential difference between two plates usually placed alongside the permanent magnets. The full Lorentz force then disperses the ions, $\mathbf{F} = q(\mathbf{E} + \mathbf{v} \times \mathbf{B})$, with the electric dispersion in the z -direction and the magnetic in the y -direction resulting in a parabolic trajectory. Thus with measurement of the fields the author wrote a particle tracking code to identify the parabolas left on the detector (typically CR-39) which are unique to the ion charge to mass ratio, Z/A .

3.2.5 Ionography

Ionography, which is more commonly known as ion radiography, uses ions as an active diagnostic tool. A beam of laser plasma produced ions is used to obtain a projected image of an object. The energy and spatial profile of the ion beam is modified by both object particle collisions and interactions with electromagnetic fields.

Ion sources can be used to obtain images of low contrast objects which are too dense to be probed by optical radiation. The broad spectrum of laser plasma accelerated ions along with a suitably filtered detector allows tuning of the ion energy used such that their mean free path is just longer than the sample thickness hence allowing high contrast images. Previously this kind of imaging required massive, expensive accelerator facilities. The last two decades have enabled ionography to be carried out on university based laser systems with relatively low intensities such as the JLITE-X setup in Japan ($10^{16} \text{ W cm}^{-2}$) which achieved 600 nm resolution with ions produced from 0.5 μm CO_2 clusters when probing spiders [96].

A unique application of ion beams, other than probing object density, is to probe transient electromagnetic fields. Here, the ion beam is typically imprinted with a mesh before passing through the sample and any distortions of the mesh image on the detector are modelled using ray tracing to simulate the fields. The advantage of laser accelerated ions is the picosecond time resolution achievable for observation of the evolution of transient fields in, for example, a laser-plasma interaction experiment [97] [98] [99].

3.3 Optical Diagnostics

In this section we shall introduce a range of optical methods for spatial and temporal resolution of plasma density. We concentrate on the techniques of shadowgraphy, dark-field schlieren imaging and interferometry. All these techniques rely on the refraction of light as the refractive index of the plasma is modified by changes in electron density, n_e . Detecting the additional phase of the probe light enables recovery of this electron density. However, above the critical density techniques such as ionography are necessary to probe further as the light is not transmitted.

3.3.1 Schlieren Imaging

The speed of light in a sample with refractive index η is given by $c_{sample} = c/\eta$. Using Huygen's principle: 'every point on a propagating wave front serves as the source of spherical secondary wavelets, such that the wave-front at a later time is the envelope of these wavelet' [100] we can visualise some of these wavelets slowing down as they pass through a region of different density and refractive index and thus the overall wave-front envelope tilting - known as refraction. More rigorously we can generalise this to a deflection from the original propagation direction along the x-axis of θ caused by a variation in the three-dimensional refractive index $\partial\eta$ [101]:

$$\theta_y = \frac{1}{\eta_0} \int \frac{\partial\eta}{\partial y} dx \quad (3.9)$$

$$\theta_z = \frac{1}{\eta_0} \int \frac{\partial\eta}{\partial z} dx \quad (3.10)$$

where η_0 is the reference refractive index such as that of the homogeneous surrounding medium. Clearly these deflections, θ , will produce variations of intensity in an initially uniform probe. Projecting these variations onto a detector is known as *shadowgraphy*. The detected intensity variation, ΔI , from the initial intensity, I_0 , a distance L from the sample is therefore given by:

$$\Delta I = I_0 L \left(\frac{\partial^2}{\partial y^2} + \frac{\partial^2}{\partial z^2} \right) \int \eta dx \quad (3.11)$$

where a uniform density gradient deflects the image rays uniformly such that the increase in intensity with respect to the initial intensity is proportional to $\nabla^2\eta$ - shadowgraphy thus diagnoses the second spatial derivative of the refractive index.

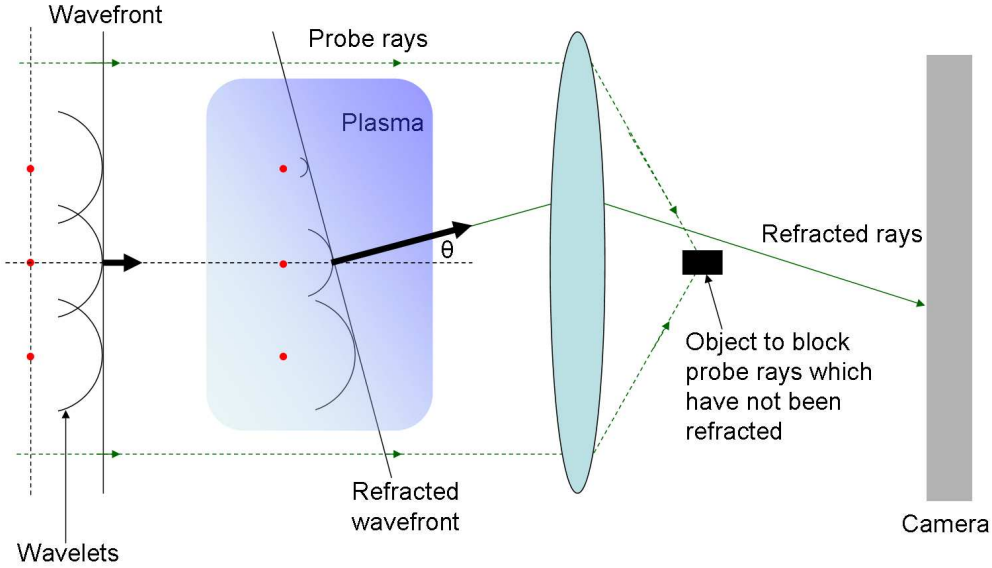


Fig. 3.18: Schematic of dark-field schlieren imaging. Probe rays are refracted as stated by the Huygens principle as they travel through regions of inhomogeneous electron density which exhibit refractive index gradients. There refracted rays no longer focus at the focal distance f of the lens. The original undeviated rays can be blocked allowing an image of the phase object to be obtained on the detector.

Imaging the sample with a lens maps every point in the sample plane to a unique point in the image plane - losing the angular deflection information described above. However, the refracted rays will come to a different focus than the initial undeviated rays. If a small object, such as a $100\ \mu\text{m}$ needle is placed at the focus of the initial probe beam only the refracted rays will reach the detector plane. Hence the intensity observed is a function of the deviation, or the refractive index gradient. This extension of shadowgraphy is known as *dark-field schlieren* and is shown schematically in Fig. 3.18.

3.3.2 Streaked Schlieren Imaging

Schlieren imaging allows two dimensional spatial resolution of the plasma density gradient. Most of the systems we wish to study evolve rapidly over a nanosecond, or sub-nanosecond, time scale necessitating the addition of time resolution. Alistair Moore *et al.* developed the *streaked schlieren* technique which allows time resolution at the expense of a spatial dimension [102].

A *streak camera* as shown in Fig. 3.19 relays photons from the input slit onto a photocathode where they are converted to electrons via the photoelectric effect. The electrons are then collimated with an electrostatic Einzel lens and accelerated with an anode field. The propagating electron experience a time varying perpendicular electric

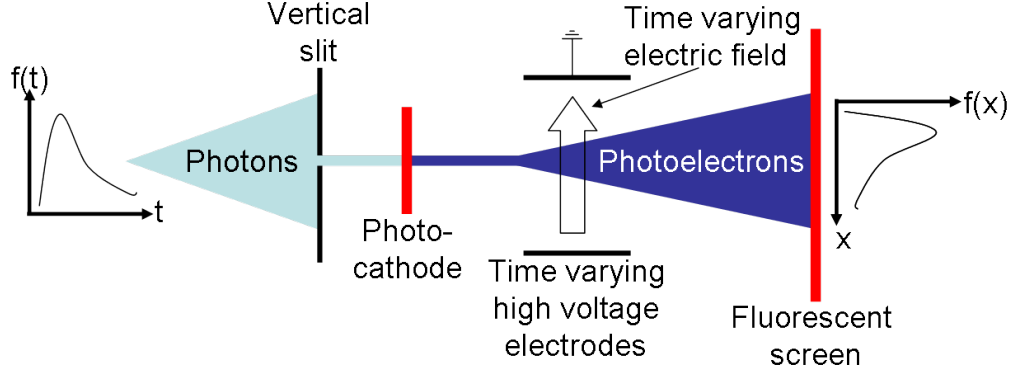


Fig. 3.19: Plan view schematic of a streak camera. The incident photon signal is converted to electrons via a photocathode and then ‘streaked’ with a time varying high voltage electric field in order to allow a camera to capture an image of one spatial dimension against time from a fluorescent screen.

field that maps them to a different height on the multi-channel plate (MCP) detector according to the time at which they see the electric field. Knowledge of the time varying electric field allows the image obtained with camera and a phosphor screen behind the MCP to be converted to a temporal intensity profile with picosecond resolution. Typically a long pulse laser ($\sim$ nanosecond) is used to probe the object for the duration of the streak. We also make use of an x-ray streak camera which operates along a similar principle with an ‘open’ photocathode for probing to higher object densities in the warm dense matter experiments described in later chapters.

3.3.3 Interferometry

Shadowgraphy and schlieren images produce clear and typically easy to understand images of density gradients. However, to quantitatively analyse the electron density interferometry is necessary. Interferometry overlaps a reference light beam which has travelled through neutral vacuum and an object beam which acquires a relative phase shift from the local refractive index, $\Delta\eta$, of the sample. The resulting interference pattern can potentially be deconvoluted to obtain the phase shift and hence the refractive index which for a known material allows approximate recovery of n_e . The electron density is only approximate as we make the assumption that the refractive index is dominated by electron effects whereas ion, neutral gas and the Kerr effect also contribute. Under this assumption [103]:

$$\eta \approx \eta_e = \sqrt{1 - \left(\frac{\omega_{pe}}{\omega}\right)^2}. \quad (3.12)$$

For $n_e \ll n_{crit}$ this is simplified to $\eta \approx 1 - \frac{n_e}{2n_{crit}}$.

The difference in refractive index from vacuum is given by $\Delta\eta \approx \eta_{plasma} - 1$ then the total phase shift, ϕ , is [103]:

$$\phi = \frac{2\pi}{\lambda} \int_{path} (\eta_{plasma} - 1) dx \quad (3.13)$$

where λ is the light's wavelength and *path* denotes the integral is over the trajectory through the entire plasma. Substituting in our value for the refractive index above the phase shift becomes:

$$\phi = -\frac{\pi}{\lambda n_{crit}} \int_{path} n_e dx. \quad (3.14)$$

Fig. 3.20 and Fig. 3.21 show the two types of interferometer used in this thesis, the *Mach-Zehnder* and the *Michelson*. In both cases the sample is imaged with a lens in order to remove the effects of refraction described above in the shadowgraphy technique.

Both styles of interferometer produce similar interference results however they differ in the method they use to split and recombine the beams. The Michelson was originally used for single blast wave experiments as it requires fewer optics, particularly important within the limited space of a target chamber. Here, a large beam is used where half passes through the sample and the other half acts as the reference. A rooftop prism inverts fifty percent of the beam after a beam splitter such that when they are recombined fringes can be obtained. The prism is usually mounted on a translation stage to match the path length of both arms after the beam-splitter in order to obtain high contrast fringes.

For later experiments involving colliding perpendicular blast waves the Michelson system was problematic as the reference half of the beam always captured some of the blast wave. To overcome this a Mach-Zehnder interferometer was used. Here, the reference beam and probe beam are split with a beam-splitter before the sample and recombined afterwards with another beam-splitter. Again, one of the beam-splitters is mounted on a translation stage so that the optical path lengths of the two arms can be matched to within a pulse length ($\sim 150 \mu m$).

In both cases the fringes can be mathematically described by [104]:

$$g(x, y) = a(x, y) + b(x, y) \cos(2\pi k_0 x + \phi(x, y)) \quad (3.15)$$

where a and b are scale factors arising from intensity variations, k_0 is the unshifted spatial

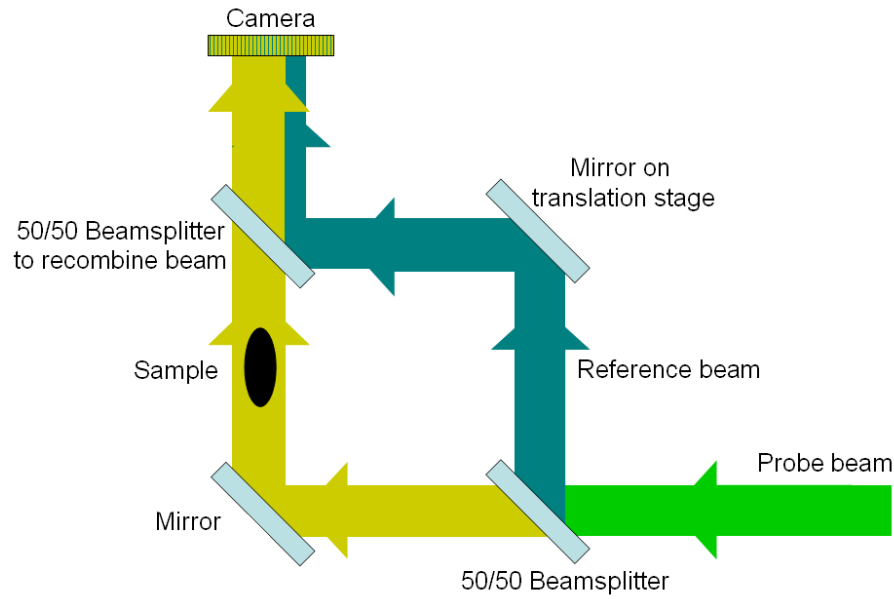


Fig. 3.20: Schematic of a Mach-Zehnder interferometer. Here, the original 532 nm probe beam is split via a 50/50 beam splitter into a reference arm and a sample arm. The phase object being probed can be as large as the sample beam, an advantage over the Michelson style. The beams are then recombined via a beam-splitter and matched in path length using a translation stage to achieve high contrast fringes on a camera.

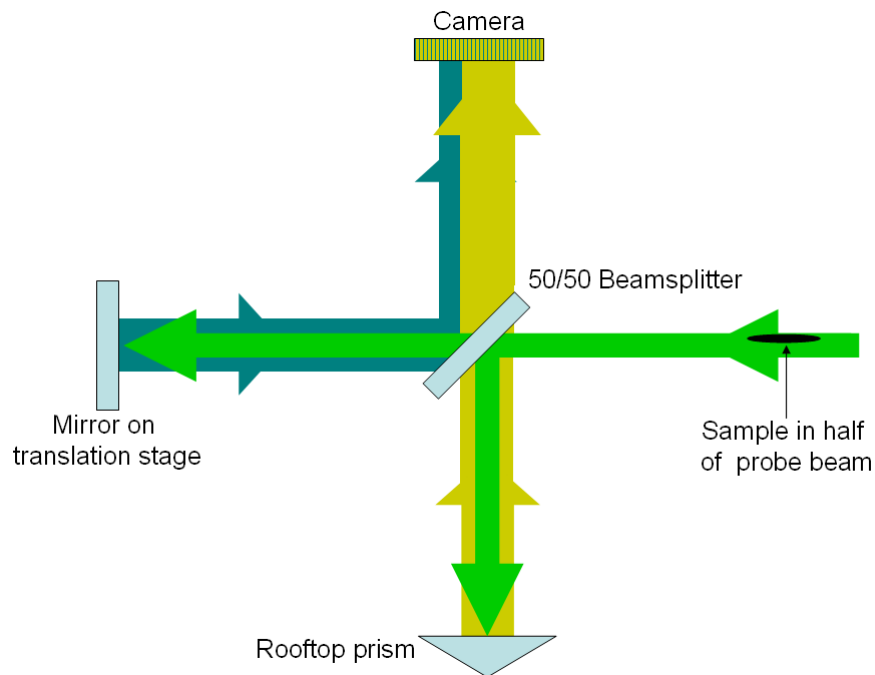


Fig. 3.21: Schematic of a Michelson interferometer. Here, only one beam is used within the chamber allowing for fewer optics. However, the sample can only occupy one half of the incident probe beam as the other half is used as the reference beam. This is achieved after beam splitting with a 50/50 beam splitter and inverting with a rooftop prism such that one the beams are combined the reference half overlaps the sample half allowing, with path length matching, phase information recovery from fringes observed on a camera.

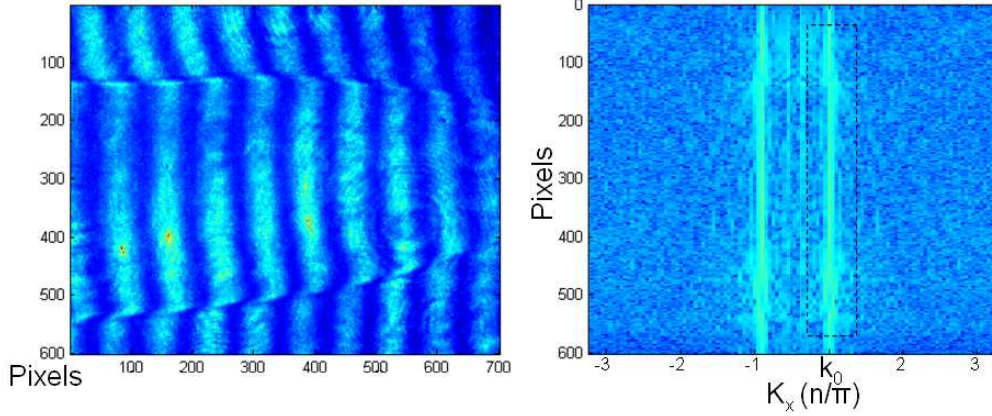


Fig. 3.22: Recovering phase information from interferometry. The left image shows fringes recovered from a Michelson interferometer of a horizontal cylindrical blast wave in argon from the Glass laser. The right image shows the modulus of the 1D Fourier transform acquired using the PAT program. The two sidebands are clearly visible with the positive one shown in a dashed box. This is windowed and shifted by k_0 to the origin in order to carry out an inverse transform and obtain the phase.

fringe frequency and ϕ is the imparted phase shift which gives rise to the fringe shift.

Rewriting in exponential form to take the Fourier transform:

$$\mathcal{F}_x(g(x, y)) = \mathcal{F}_x(a(x, y)) + \mathcal{F}_x\left(\frac{1}{2}b(x, y)e^{i\phi(x, y)}\right) + \mathcal{F}_x\left(\frac{1}{2}b(x, y)e^{-i\phi(x, y)}\right) \quad (3.16)$$

where the final term is the complex conjugate of the previous term. Fig. 3.22b shows the side bands arising from these two terms which have been shifted by k_0 from the centre. If one of these sidebands is selected with a Hann window and shifted back to the centre an inverse Fourier transform can be carried out to obtain $\frac{1}{2}b(x, y)\exp i\phi(x, y)$. Taking the imaginary part of the natural logarithm thus returns the phase. This process does not take into account phase shifts of more than π , these potential discontinuities have to be ‘unwrapped’ manually. In order to recover electron density values from the phase the geometry of the sample needs to be accounted for as the phase shift is an integral of contributions from the whole path length of the sample. Fortunately with the majority of our blast waves we can assume a cylindrical geometry which simplifies our analysis by allowing the use of an *Abel inversion*. The Abel inversion allows us to recover an axially symmetric function from a projection. The inverse radial transform in cylindrical coordinates is [103]:

$$n_e(r) = \frac{\lambda}{\pi^2 n_{crit}} \int_r^R \frac{d\phi}{dy} \frac{dy}{\sqrt{y^2 - r^2}} \quad (3.17)$$

where R is the outer radius of the blast wave where $n_e = 0$ - typically removal of back-

ground data is carried out with a null shot.

For the majority of interferometry analysis contained in this thesis the PAT (Phase calculation Abel inversion and Tomography program) software developed by Matthias Hohenberger and James Lazarus which incorporates much of the above methodology automatically via a graphical user interface is used.

3.4 Cluster Production

As discussed in the *different target media* section, in order to produce blast waves experimentally we make use of clusters as a target to absorb the laser energy. These clumps of thousands of atoms held together by van der Waals forces or hydrogen bonding can be formed using a gas nozzle. As the gas, held at a high backing pressure (typically $\sim 30 \text{ bar}$), expands adiabatically through a nozzle in to vacuum the atoms cool and form clusters. The empirical Hagen parameter, Γ^* , relates the initial conditions to the onset of clustering [105]:

$$\Gamma^* = k \frac{(d/\tan \alpha)^{0.85}}{T_0^{2.29}} P_0 \quad (3.18)$$

where d is the diameter of the nozzle (μm), α the expansion half angle ($\alpha = 45^\circ$ for sonic nozzles, $\alpha < 45^\circ$ for supersonic), P_0 the backing pressure (mbar), T_0 the pre-expansion temperature (Kelvin), and k a constant related to bond formation ($k = 1650$ for Ar, 184 for H, 2890 for Kr and 5500 for Xe gas [106] [107]).

For values of $\Gamma^* > 100$ clustering can typically be observed. A scaling for the average number of atoms per cluster, $N = 33(\Gamma^*/1000)^{2.35}$, was verified in the regime $1000 < \Gamma^* < 7300$ by Hagen [108].

This thesis makes use of a modified series 99 solenoidal valve nozzle from Parker Hannifin General Valve corporation shown in Fig. 3.23. The modifications involve elongating the nozzle protrusion to allow for multiple laser beams at different angles and also encasing the nozzle in a copper cooling jacket to control T_0 . A 25 V Parker Hannifin Iota One driver pulsed the solenoid allowing gas to pass through the $500 \mu\text{m}$, 45° nozzle.

For the purposes of calculating interaction volumes, ideal target height below the nozzle and ion transport through the cluster media it is useful to calculate the density as a function of height below the nozzle. Miller details the free isentropic ideal gas expansion

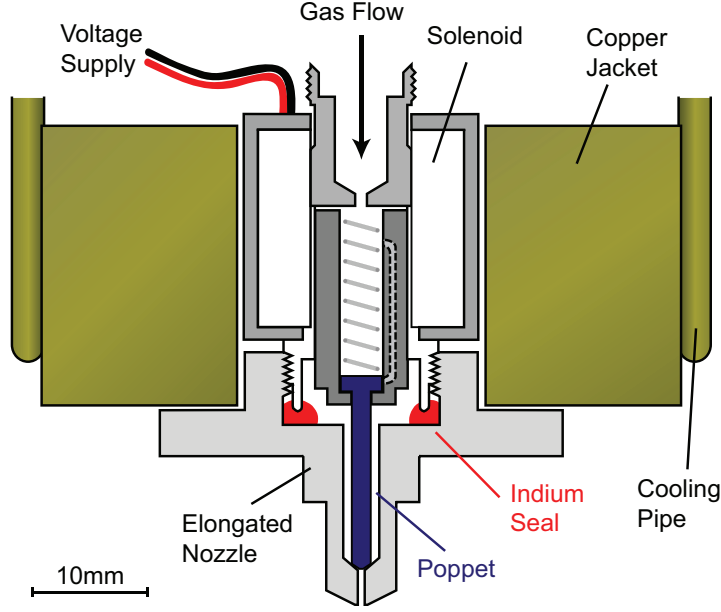


Fig. 3.23: Schematic of modified cryogenic gas nozzle. The series 99 valve was modified by reducing the circular ending nozzle to a rectangle with a width of the solenoid and also extending the nozzle opening with an elongated poppet made from diamond solution. A copper jacket with nitrogen gas cooled by liquid nitrogen flowing through it was used to cryogenically cool the gas. Figure courtesy of Dr Matthias Hohenberger.

[109] which leads us to the centreline average density:

$$n = n_0 \left(1 + \frac{\gamma - 1}{2} M^2 \right)^{\frac{1}{\gamma - 1}} \quad (3.19)$$

where n_0 is the backing gas density, γ is the gas polytropic index (typically 5/3). M , the Mach number, is given by:

$$M = A \left(\frac{x}{d} - B \right)^{\gamma - 1} - \frac{\gamma + 1}{2(\gamma - 1)A \left(\frac{x}{d} - B \right)^{\gamma - 1}}. \quad (3.20)$$

A and B are constants which weakly vary with γ - $A = 3.26$ and $B = 0.075$ for $\gamma = 5/3$. These equations are only valid for $\frac{x}{d} > 2.5$.

Typical experimental backing pressures of ~ 40 bar of argon gas using the modified extended nozzle, which allows for a greater diagnostic probing solid angle, with a $500 \mu\text{m}$ aperture at room temperature results in clusters of radius $\sim 35 \text{ nm}$ with a local density of $\sim 10^{22} \text{ cm}^{-3}$ in the cluster and an average atomic density of $\sim 10^{19} \text{ cm}^{-3}$ at a distance of 2.5 mm below the nozzle.

3.5 Summary

In this chapter we begin with a brief description of the fundamentals of short pulse, high intensity lasers which are used in all our experiments. We particularly concentrated on the breakthroughs of Q switching, mode-locking, chirped pulse amplification and optical parametric amplification which are necessary to fully describe the operation of the Imperial College Nd:Glass and Vulcan lasers used in the majority of this thesis. Then we introduced ion diagnostics with particular reference to the first absolute calibration of radiochromic film and a home built magnetic spectrometer to low energy protons (0.07 - 4 MeV). Optical diagnostics and methods for the extraction of expansion velocity and electron density were subsequently described. Finally, the gas nozzle used for the production of clusters and the associated characteristics were established.

Chapter 4

Proton Heating Towards Warm Dense Matter

This chapter describes warm dense matter generation and characterisation experiments based on the theory and experimental methods described in chapters 2 and 3. The experiments described here were conducted in 2008 on the Helen laser at AWE and 2010 on the Vulcan (Target Area West) laser at the Rutherford Appleton Laboratories. Both the experiments were conceived of and run by the Plasma Physics Department at AWE. My role, which is the main topic of this chapter, is determination and analysis of the ion spectrum which heats the WDM, and associated deposited energy calculations. Also included are results from spatially and temporally resolved density analysis and preliminary ultra-violet emission results.

In this chapter improvements are detailed for accurate measurements of the low energy ion spectrum using a combination of a well calibrated magnetic spectrometer, Thomson parabola and radiochromic film stacks. Previous methods relying solely on radiochromic film overestimate the lower end of the ion spectrum by more than a factor of two. Discrepancies between different methods of calculating energy deposition highlight the errors in codes which use cold stopping powers which may not be accurate for warm, partially ionised matter. Finally, attention is drawn to the long time of flight dispersion of the ion beam which results from a higher than expected proportion of carbon and oxygen ions, which will cause non-uniform, non-isochoric heating.

4.1 Measurements Required for Equation of State Determination

The challenge of generating microscopically large, near-uniform samples of WDM in the laboratory is significant and, as described in previous chapters, several routes are being pursued by a number of research groups worldwide. However, once obtained, sufficiently accurate diagnosis of this state of matter is an equally significant challenge. Standard optical techniques cannot penetrate these densities due to their high opacity on account of the high plasma frequencies. Also, many standard spectroscopy techniques are not applicable at these low relatively temperatures. Very high spatial and temporal resolution is also necessary to resolve the sample size as the state of matter evolves.

The key underlying measurements required to help resolve the equation of state are sample temperature, T , and density, ρ , observed simultaneously over a uniform volume as they evolve. Observations of the energy deposited in the sample allow us to ensure isochoric heating and comparisons with temperature show heating and loss mechanisms within the sample. In determining the equation of state we also need to clarify the sample's ionisation state over time. Deposited energy will excite, ionise and increase the average kinetic energy of the sample. Current models disagree to a significant extent over the ionisation states in the temperature range we investigate [110]. Knowledge of the ionisation state and population of atomic levels is critical for derivation of the equation of state and opacities as this will measure the relative electron to ion contribution to pressure.

4.2 Experimental Setup for Heating and Probing

The method we use for heating and probing warm dense matter was developed over several experiments at AWE, Aldermaston by Dr David Hoarty and co-workers. In particular two experiments provide data detailed in this thesis - on the Helen laser in 2008 and the Vulcan TAW laser in 2010. Both these investigations use a similar experimental setup however the latter adds several diagnostics. The initial Helen experiment was used as a proof of principle platform to streak XUV images of an ion heated sample. Although RCF stacks were used in 2008, no characterisation of the low energy (< 4 MeV) or heavy ion spectrum ($Z > 1$) was carried out.

The 2010 experiment developed the XUV imaging to higher resolution and added UV streaked imaging, ion spectroscopy and low energy proton detection. Fig. 4.1 shows the

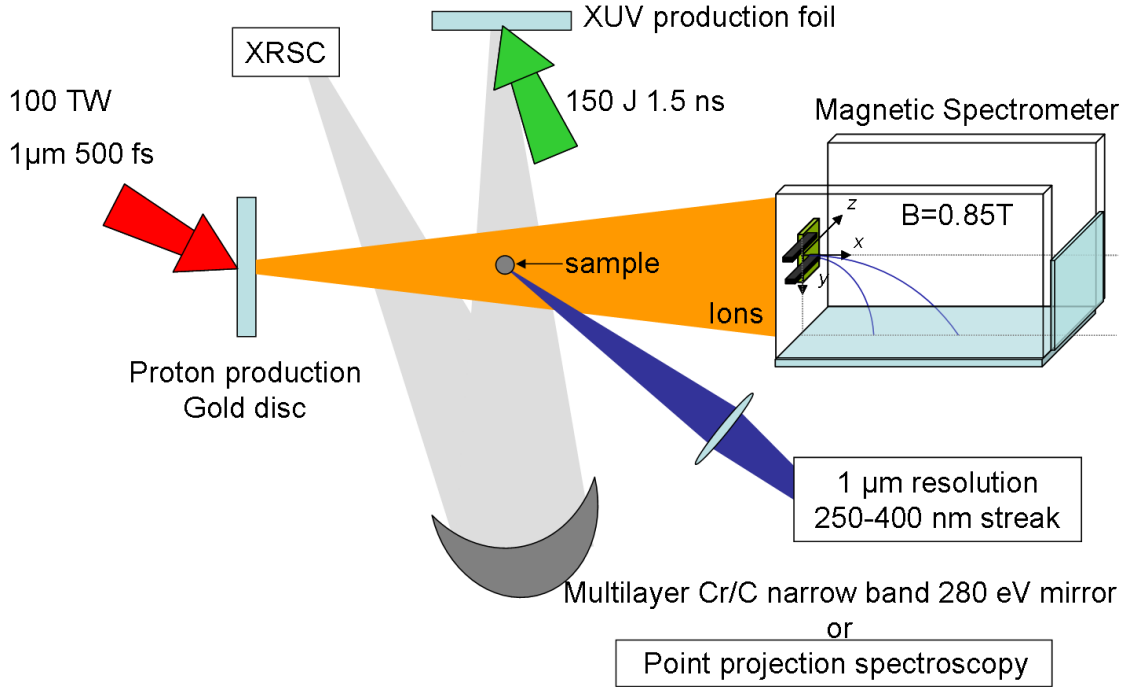


Fig. 4.1: Top view schematic of the WDM Vulcan TAW experiment 2010. A 100 TW short pulse CPA beam incident on a gold primary target produced an ion beam which heated our sample. The ion beam spectrum was sampled with a magnetic spectrometer. For some shots the spectrometer was replaced, or used in conjunction with, a Thomson parabola or a RCF stack. A second long pulse green beam was used to produce XUV light via a foil interaction. This XUV imaged the sample evolution via a narrow band multilayer mirror onto an X-ray streak camera (XRSC). An achromatic lens was used to streak the UV self-emission from the sample. Towards the end of the experiment the multilayer mirror was replaced with a crystal spectrometer for point projection spectroscopy.

experimental setup with these additional diagnostics.

A 100 TW short pulse CPA beam incident at an angle of $\sim 22^\circ$ onto a 200 μm diameter, 15 μm thick gold disc (primary target) was used to produce an ion beam via the target normal sheath acceleration process. The majority of electrons and x-rays produced exit the rear side of the disc along the laser axis - the geometry was designed such that they do not heat the target. The distance to the sample (secondary target) was set at 250 μm . Several types of metal wire were heated as detailed in Table 4.1. Aluminium was used for the majority of shots as it is a commonly investigated metal with much literature focussing on it and will act as a suitable benchmark. The other materials are of interest for work such as that on the National Ignition Facility.

After the ion beam passes through the sample the spectrum is measured using the magnetic spectrometer (§3.2.3) positioned 300 mm from the sample. This was used in

Material	Wire Diameter, μm
Aluminium	10
Aluminium	5
Carbon	7
Gold	10
Copper	10
Copper	20

Table 4.1: List of the materials and sizes probed in the warm dense matter regime. All samples had the form of wires and were prepared by either the AWE or Rutherford laboratories target preparation facilities.

conjunction with an RCF stack to look at the higher energy protons (>2 MeV), spatial distribution and absolute fluence. For some shots the spectrometer was replaced with a Thomson parabola in order to recover the ion species present in the heating beam.

Up to two green long pulse beams, each with 150 J, 1.5 ns in a 500 μm focal spot were used to produce XUV radiation from a 50 μm gold foil. The XUV radiation after passing through the sample wire was focused with a curved multilayer sacrificial mirror through a large throw (~ 5 m) to give $\times 50$ magnification. The Cr/C mirror was narrow band, only reflecting 280 eV photons. The image was captured using an AWE home-built x-ray streak camera capable of ~ 10 ps time resolution. Fine meshes were used in the object plane to determine the back lighter spatial resolution of 2 μm .

For the latter shots in the 2010 experiment the multilayer mirror was replaced with a crystal spectrometer from Belfast university. Here M-band x-rays were generated with a Hafnium foil which bracket the aluminium K-shell transition arrays in order to measure ionisation via point projection spectroscopy. This diagnostic had limited success.

The final diagnostic imaged the ultra-violet (250-400 nm) self-emission from the wire using an achromatic lens with a long throw to obtain $\times 50$ magnification and theoretically a close to diffraction limited resolution of 1 μm . The image was captured on a Sydor streak camera which allowed ~ 10 ps time resolution.

4.3 Results

In this section the results and analysis are presented from the setup described above for the 2008 and the 2010 experiment. Mainly the results of the deposited energy through ion spectrum analysis are covered. However, also briefly shown are the results of density analysis via XUV streaking by AWE for equation of state recovery and UV streaking analysis by M. Hill at AWE [111] as an aid to determine the uniformity of ion heating.

4.3.1 Time Resolved Density

The initial proof of principle Helen 2008 experiment relied solely on the XUV streaking diagnostic. Fig. 4.2a shows the raw data from one such shot with 57 J, 0.5 ps of laser energy focused onto a 15 μm gold disc, and a test sample of 10 μm aluminium wire. Fig. 4.2b show the results from NYM rad-hydro [112] simulations carried out by AWE which allow us to estimate the energy deposited and temperature reached based on the rate of expansion. These simulations however have a huge potential for error as we have already shown that models of this type cannot accurately predict the behaviour of WDM without significant further work. Also, here we model the desirable *instantaneous* energy deposition which is clearly not the case without a purely mono-energetic proton spectrum. Lastly, we assume an initial temperature from the deposited energy, however as the state of the sample evolves unknown energy loss mechanisms require a direct measurement of the temperature.

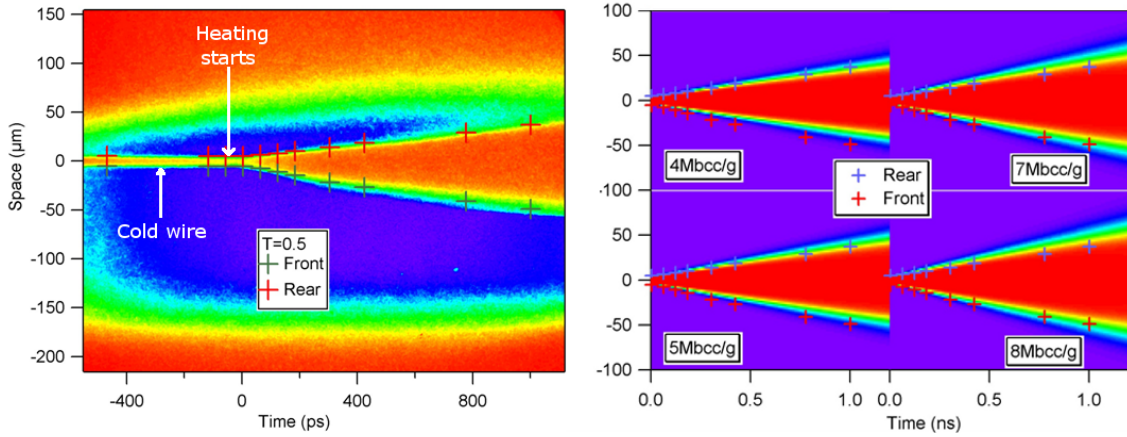


Fig. 4.2: a) Raw XUV radiography streak camera data taken on a 57 J, 0.5 ps Helen laser shot. The sample is a 10 μm aluminium wire which can be seen at cold, solid density for $t < 0$ and at zero on the y-axis. As the heating protons arrive at $t > 0$ the wire is seen to expand into vacuum beyond the initial 10 μm diameter. Contours were established, as shown by the crosses, at 50% transmission of the XUV backlighter - the gradient of these contours is the surface expansion velocity. The front and rear show near uniform expansion at a rate of $\sim 4 \times 10^4 \text{ ms}^{-1}$. b) NYM Rad-hydro [112] simulations carried out by AWE with uniform energy deposition followed by free expansion of a 10 μm aluminium wire. Varying the initial energy deposition, as shown in these four subplots, allows comparison of the rate of expansion to the experimental data. The simulations suggest initial energy deposition of $4 \times 10^5 \text{ Jg}^{-1}$ ($\sim 14 \text{ eV}$ equivalent temperature) on the rear and $8 \times 10^5 \text{ Jg}^{-1}$ ($\sim 20 \text{ eV}$) at the front surface.

In order to recover the density from the transmission we can use Beer-Lambert's law:

$$\rho(z) = -\frac{\ln(1 - A(z))}{\kappa \Delta x} \quad (4.1)$$

where A is the absorption, κ the opacity and Δx the thickness at a point on the wire, z . Values for opacity are taken from opacity tables. In our density regime the opacity is dominated by bound-free absorption and as such does not vary strongly with density and temperature unless ionisation occurs into a new shell - requiring significant additional energy for instance into the l shell in aluminium [113] [9]. Finally an Abel inversion needs to be carried out to recover the cylindrical geometry for radius r from the path integrated value:

$$\rho(r) = \frac{1}{\kappa(r)\pi} \int_r^{r_{max}} \frac{d}{dz}(-A(z)) \frac{dz}{\sqrt{z^2 - r^2}}. \quad (4.2)$$

Is it possible to use this measured density to calculate pressure without explicit temperature and ionisation equation of state measurements? Foord *et al.* [114] describe a method which allows us to obtain an estimate of pressure, P , from single values of density within the limits of an isochoric, self-similar, isentropic planar system.

Starting with the equations of continuity and motion (§2.1.5):

$$\frac{\partial p}{\partial t} + \frac{\partial(\rho v)}{\partial x} = 0 \quad (4.3)$$

$$\frac{\partial v}{\partial t} + v \frac{\partial v}{\partial x} = -\rho^{-1} \frac{\partial p}{\partial x} \quad (4.4)$$

where p is pressure, ρ is mass density and v is plasma velocity. Isentropic flow implies $\frac{\partial p}{\partial x} = \frac{\partial p}{\partial \rho} \frac{\partial \rho}{\partial x} \equiv C_s^2 \frac{\partial \rho}{\partial x}$ where C_s is the sound speed. This allows us to rewrite Eq. (4.4) for an adiabatic expansion [114]:

$$\frac{\partial v}{\partial t} + v \frac{\partial v}{\partial x} = -C_s^2 \rho^{-1} \frac{\partial \rho}{\partial x}. \quad (4.5)$$

Equations Eq. (4.3) and Eq. (4.5) can then be solved via substitution [114] to yield $\frac{\partial(\rho C_s)}{\partial x} = -\frac{\rho}{t}$. Thus, integrating the density across the sample yields the sound speed:

$$C_s = -(\rho t)^{-1} \int_{x_R}^x \rho dx + \frac{\rho_0 C_{S0}}{\rho} \approx (\rho t)^{-1} \int_x^{x_{max}} \rho dx \quad (4.6)$$

where x_R is the displacement of the rarefaction at time t and ρ_0 , C_{S0} are the density and

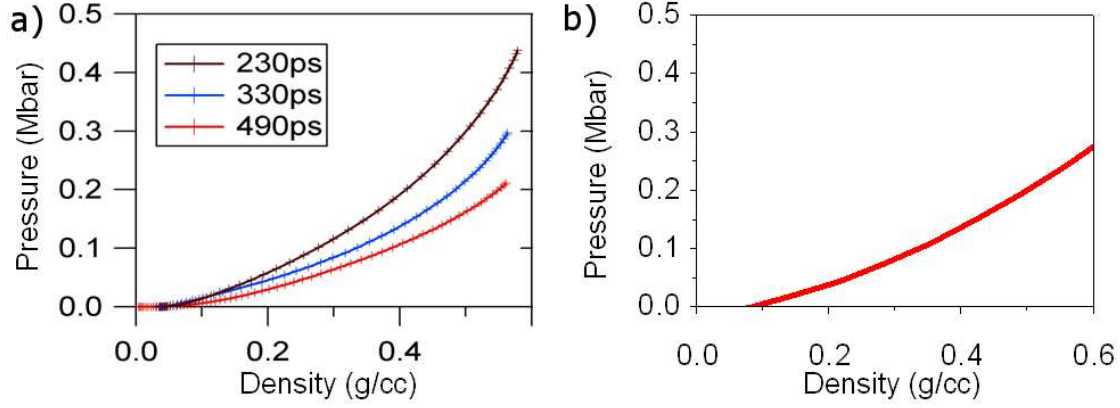


Fig. 4.3: a) Measured isentrope using the Foord method to obtain pressure from single value of density at different times in a 10 μm aluminium wire expansion. The lines at 330 and 490 ps occur at late time periods where the self-similarity assumption breaks down and have increased error. Analysis carried out by AWE. b) Theoretical isentrope adapted from Foord *et al* [114] for a 10-11 eV solid density plasma.

sound speed of the initial material. The pressure is thus the integrated C_s^2 :

$$P(\rho) = \int_{\infty}^x C_s^2 \frac{\partial \rho(t)}{\partial x} dx. \quad (4.7)$$

Thus, using the data shown in Fig. 4.2a we obtain the P, ρ curves shown in Fig. 4.3a. The curves do not overlap when taken at different times. We attribute this apparent non-isentropic behaviour to four main sources of error. Firstly, non-uniform heating affects the validity of the adiabatic limit and this is discussed later. Secondly, since our cylindrical wire does not fulfil the planar geometry assumption required for this method we underestimate the pressures. Thirdly, Foord sets boundary conditions such that the density sound speed product is zero at the sample edges. At early time this introduces a larger error than the later times assessed by Foord [113]. Lastly, at late times when the rarefaction reaches the centre of the sample it varies from the self-similar limit, and as such the Foord method cannot be applied in this central region [113]. As such we choose a time during the expansion where the Foord method is most valid away from these deviations - typically ~ 200 ps after the initial ion heating. For further discussion on the applicability of the Foord method the reader is directed towards [115]. The data broadly agrees with simulations by Foord shown in Fig. 4.3b.

The lack of penetration of the 280 eV photons through the high densities present at the beginning of the expansion do not allow extension of the isentrope to the highest pressures believed to be present (~ 10 Mbar).

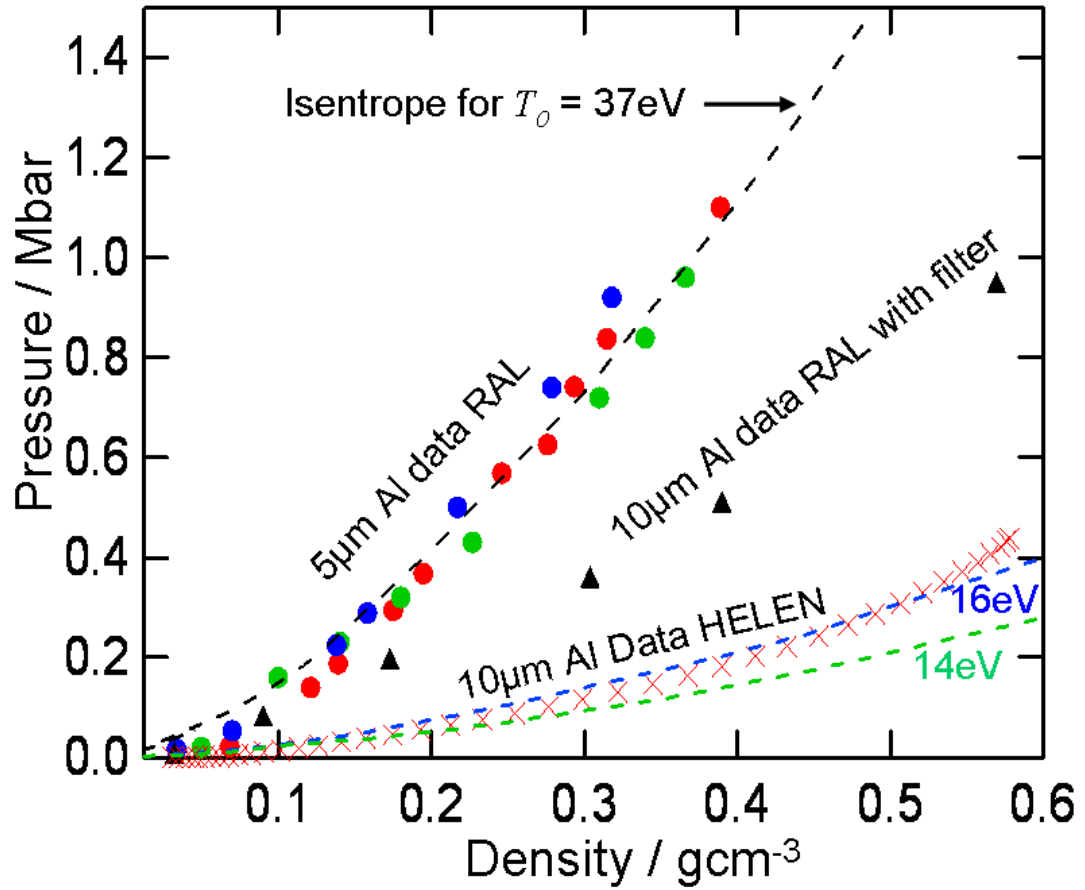


Fig. 4.4: Pressure-density plot from radiography data. The circular points were taken with the Vulcan laser using a 5 μm aluminium wire sample. The different colours represent different shots. The triangular points were also taken on the Vulcan laser using a 10 μm aluminium wire sample with a 6 μm aluminium foil between the proton source and sample in an attempt to filter out high Z ions and slow protons to provide more uniform heating. The cross points were taken on the Helen laser using a 10 μm aluminium wire sample. All densities were taken from the Abel inverted radiography and the pressures calculated using the Foord method [114]. The dashed lines are isentropes for a given initial temperature taken from the SESAME library data. Adapted courtesy of AWE.

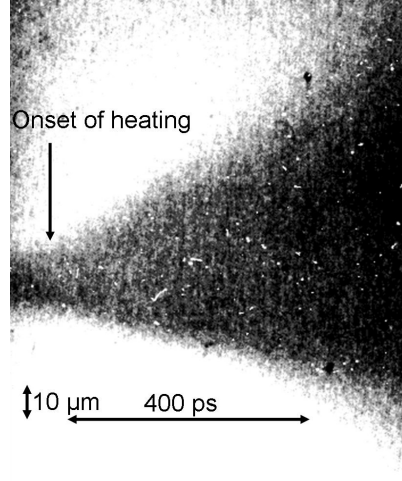


Fig. 4.5: Raw streaked XUV backlit data from 2010 experimental data. Shot 54 with 74.9 J of pre-compressor laser energy and a 10 μm aluminium sample. This shot also had a 6 μm aluminium foil to act as a heavy ion filter in order to increase the uniformity of heating. Data shows more front to back asymmetry than the 2008 data. Adapted courtesy of AWE.

The 2008 experiment, using the method and results described above, clearly showed that warm dense matter on the order of ~ 15 eV is achievable using proton heating. However, it also clearly showed that additional diagnostics are necessary to accurately recover the equation of state.

A similar analysis was carried out on the 2010 experimental data as shown in Fig. 4.4 suggesting equivalent temperatures of ~ 37 eV are observed [115]. The uniformity of expansion can be matched to the energy deposition measured with the magnetic spectrometer. Fig. 4.5 shows the processed data from a shot where a 6 μm aluminium foil was placed between the primary and secondary targets in order to remove the carbon and oxygen ions present in the ion beam which would otherwise heat the 10 μm aluminium sample very asymmetrically due to their low penetration.

4.3.2 Deposited Energy Determination from Ion Spectra

Physical models of the proton energy distribution show an approximately Maxwellian form [50]:

$$\frac{dN}{dE} = \frac{N_0}{E} \exp\left(-\frac{E}{k_B T}\right) \quad (4.8)$$

where N_0 is the initial number of protons at the initial temperature, T , of the target normal sheath acceleration thermal plasma resulting in N protons with energies E . Using

RCF stack data from both the 2008 and 2010 experiments we see an approximate match to this predicted Maxwellian distribution as shown in Fig. 4.6a. Here the RCF measurement is taken from the ion beam's centre - the same area which heats the sample. The angular distribution of energies across the ion beam is discussed in chapter Ch. 6. Using a 15° expansion half angle for the beam this translates to 1.5 ± 0.4 J of proton energy in the beam, or a $\sim 6\%$ conversion efficiency for the 2010 data.

In order to obtain deposited energy in a $10\text{ }\mu\text{m}$ aluminium wire we combine our measured spectrum with SRIM calculations for energy loss:

$$\Delta E_{dep} = - \int_0^z \int_0^E E(z, E_0) \frac{dN}{dE_0} dE_0 dz \quad (4.9)$$

where the deposited energy, E_{dep} , is the integral of our spectrum for all initial energies, E_0 , crossed with the energy remaining after distance, z , for each initial energy. In order to convert the data to useful units the volume covered by 1 Sr at the $250\text{ }\mu\text{m}$ target separation is used with the cold aluminium density to recover deposited energy in units of J/g.

We see clearly non-uniform heating as shown in Fig. 4.6b. This does not agree with the near uniform heating observed in the XUV streaking data. We attribute this to a peak in the proton distribution below the detection threshold of the RCF stack ($< 4\text{ MeV}$) [116]. To resolve this we use the calibrated low energy magnetic spectrometer as shown in Fig. 4.7a and an improved deposited energy in Fig. 4.7b. The low energy peak seen in this data clearly differs significantly from the spectrum suggested by the RCF and predicted by Eq. (4.8) - which is also based on RCF data.

Although the magnetic spectrometer gives a very accurate relative spectral shape it does not give absolute fluence levels. One method of obtaining this is to integrate the spectral distribution and divide by the solid angle of the slit. However, this method ignores the dynamic range of the detector, providing an underestimate of total fluence. As such the spectrometer spectrum is matched to that of the RCF stack which measures absolute fluence. Thus the data in Fig. 4.7 shows a combination of the RCF and magnetic spectrometer data which overlap at $\sim 4\text{ MeV}$. Using a 15° expansion half angle for the beam this translates to 0.5 ± 0.4 J of proton energy in the beam, or a $\sim 3\%$ conversion efficiency. The large error arises from the energy dependence in the angular distribution of protons, which is covered in detail in chapter 6, and also from the uncertainty in the RCF absolute fluence calibration.

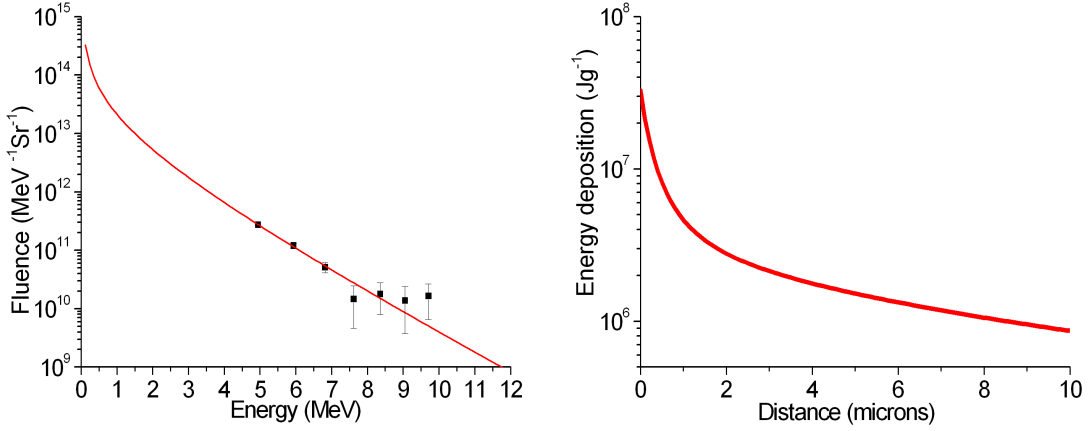


Fig. 4.6: Left: Maxwellian fit applied to fluences recovered from the RCF stack on shot 2, 79 J pre-compressor laser energy placed 65 mm from source. Only layers 3-9 provide data points as beyond this the film is beyond the dynamic range of the scanner. Right: Energy deposition obtained by combining the RCF stack spectrum with SRIM energy loss data for protons in aluminium. Over an order of magnitude difference in heating is observed between the front and rear surface of a 6 μm target wire.

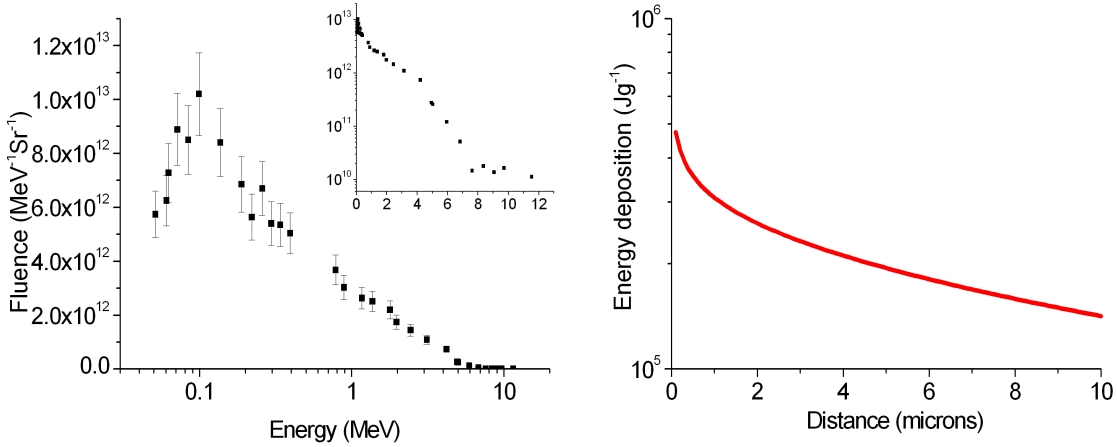


Fig. 4.7: Left: Proton spectrum captured by magnetic spectrometer on shot 56, 70 J pre-compressor laser energy. Inset shows a log-log scale for ease of comparison. The advantage of the spectrometer of RCF can clearly be seen as it extends to a lower energy. Right: Energy deposition obtained by combining the spectrum with SRIM energy loss data for protons in aluminium. The uniformity of heating is much improved on the RCF obtained spectrum and agrees with XUV data.

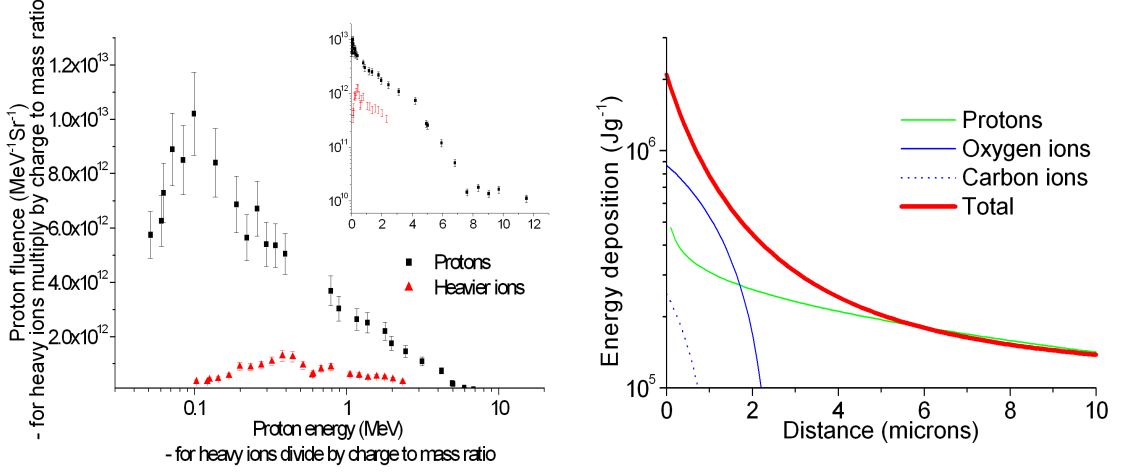


Fig. 4.8: Left: Proton and heavy ion spectrum captured by magnetic spectrometer on shot 56, 70 J pre-compressor laser energy. Inset shows a log-lin scale for ease of comparison. The Thomson parabola allows us to specify the two ion peaks as carbon and oxygen. Right: Energy deposition obtained by combining the spectrum with SRIM energy loss data for protons, carbon and oxygen in aluminium. The ions distort the heating leading to an order of magnitude difference between the front and rear surface. This does not agree with XUV data.

This more accurate low energy spectrum shows a more uniform heating at the front and rear surface of the 6 μm aluminium wire target which agrees with the density measurements. However, in the course of developing the CR-39 from the magnetic spectrometer large heavier ion pits were detected. Fig. 4.8a shows the relative proportion of protons to heavier ions. This data was obtained with the aid of aluminium filters of varying thickness on the detector plane which blocked the heavy ions but allowed through the protons, as seen in the difference in pit sizes. Clearly heavy ions ($Z > 1$) were a significant part of the heating process.

In order to resolve the species of ions present a Thomson parabola was used as shown in Fig. 4.9. With a basic particle tracking code based on hall probe measurements and a near uniform electric field the author was able to fit the parabolas to ion species. The main contaminant species are carbon and oxygen ions. Literature suggests that although we detect multiple charge states of oxygen and carbon, initially the majority of these ions are fully ionised and recombination occurs over the distance to the detector - suggesting the heating sample sees fully stripped ions [117]. Fitting carbon and oxygen peaks to the data observed with the spectrometer we obtain total deposited energy Fig. 4.8b. This corresponds to a total beam energy of 0.6 ± 0.4 J. However, the extra beam energy is deposited preferentially at the front face of the wire leading us again to the non-uniform

heating case which is not in good agreement with the XUV data.

Fig. 4.9a shows dual proton tracks which we attribute to the presence of a secondary proton source. Fig. 4.9b supports this hypothesis showing a low energy line structure offset from the primary emission detected on the RCF. Work carried out by *Beg et al.* shows that the stalk which supports the proton producing target foil, or indeed any nearby object, can develop a sheath field and emit protons as a secondary source - even if not directly illuminated by the laser [118] [119].

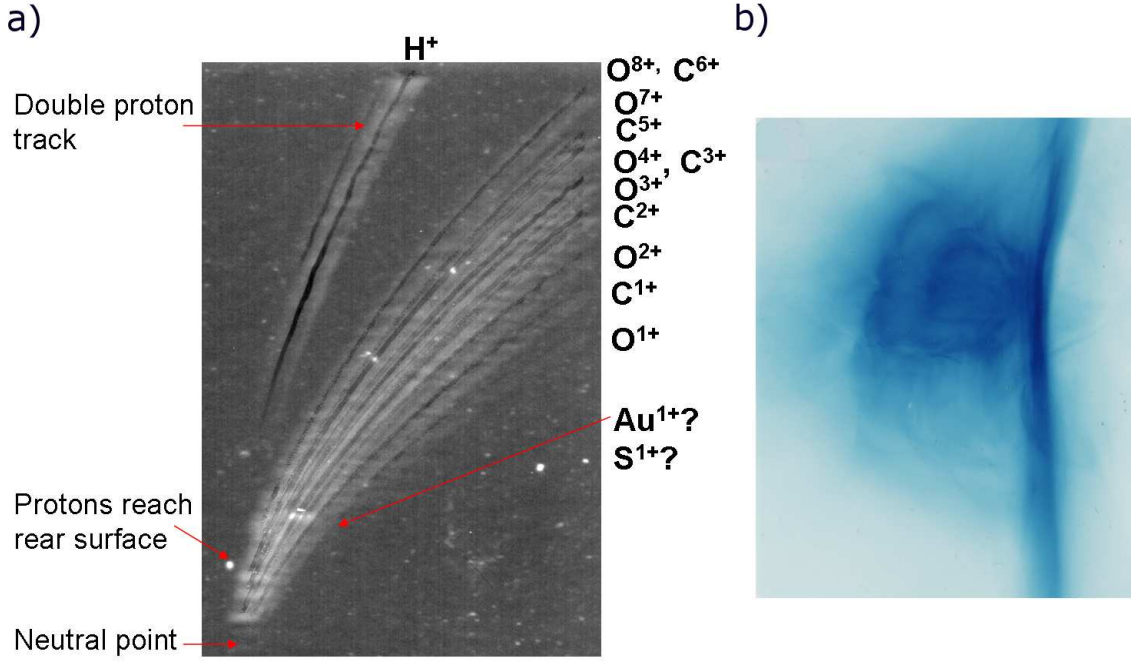


Fig. 4.9: a) Scan of Thomson parabola CR-39 from shot 46 with 84 J of pre-compressor laser energy and 7 kV potential difference between the electric plates. The parabolas are labelled using a particle tracking code. While protons, carbon and oxygen ion dominate there is faint evidence of heavier ions under the microscope which may be gold or sulphur. The proton tracks show a double line feature with a much fainter line to the left of the main parabola. We attribute this feature to two separate, displaced sources as suggested by b). b) Scan of RCF layer from a stack showing the persistent line to the right of the main ion beam attributed to a secondary source of ions arising from charging of the target stalk.

4.3.3 Uniformity of Heating from Streaked UV Emission

The UV self-emission diagnostic was partially successful, but encountered several problems. The calibrated, high time resolution (1 ps) Sydor streak camera initially intended for use failed during this run meaning calibrated colour temperature measurement were not possible and a lower time resolution device (10 ps) was used. Secondly, the high magnification ($\times 50$) used, although potentially allowing very high $\sim 1 \mu m$ spatial resolution

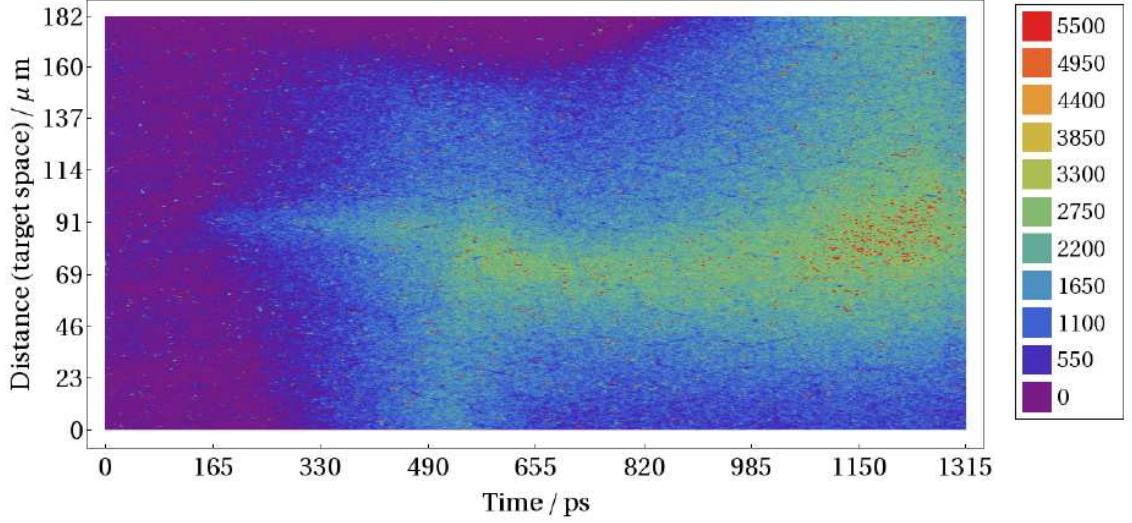


Fig. 4.10: Raw streaked UV self-emission data from shot 58 with 84 J pre-compressor laser energy. The sample was a 10 μm gold wire. The figure shows wire emission between 140 and 500 ps followed by a much brighter emission caused by collisions with the proton source target plasma as both targets expand. Figure courtesy of Dr Hill [111].

also severely affected the signal to noise ratio. The low signal levels only allowed imaging of the hotter gold wire emission as the other wires used on this run did not produce enough signal. Thirdly, the achromatic lens imaging UV light onto the streak camera was not entirely free of chromatic aberration such that different wavelengths reflected by the broadband mirror saw different effective focal lengths. This created a blurring effect on the streak camera images, reducing the spatial resolution to $\sim 5 \mu\text{m}$.

Fig. 4.10 shows the raw data obtained from the streak camera. The wire target is centred at 90 μm on this scale and the emission increases over time. At late time the primary proton beam target can be seen expanding into the heated wire plasma. Using contour analysis we can isolate this later noise from the primary target and just show the initial wire emission as seen in Fig. 4.11. As a measure of the uniformity and rapidity of heating we can take spatial and temporal averaged lineouts of this emission as seen in Fig. 4.12.

Analysis of these emission lineouts is non-trivial. The detector will only receive UV light emitted from within a skin depth of the surface of the sample. Any core emission is absorbed by this skin depth. As the wire expands and density decreases the skin depth will increase over time. Because the wire is cylindrical the line of sight skin depth to the detector and the volume of emitting gold also varies across the sample. The resolution

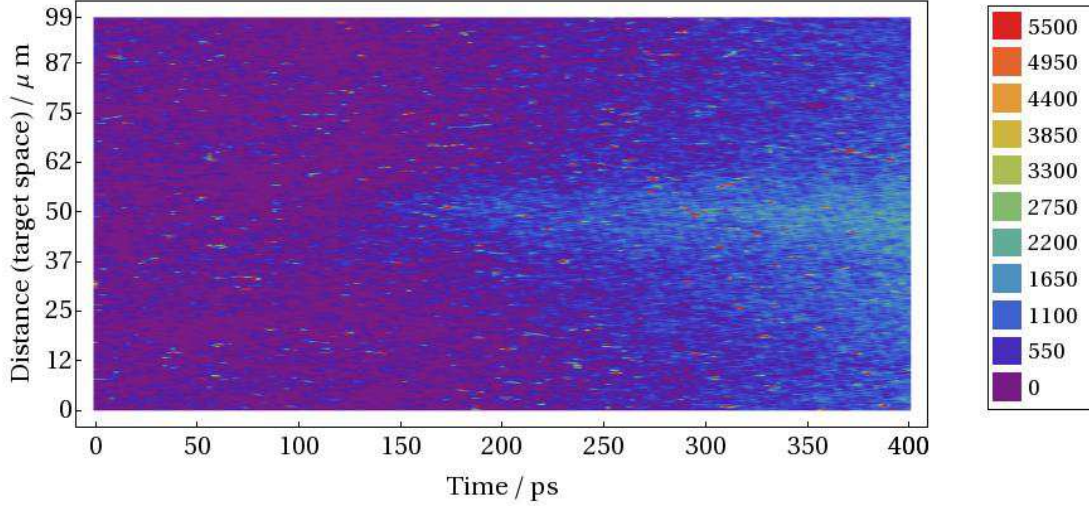


Fig. 4.11: UV emission data isolated from proton source self-emission background. Colour contour analysis was carried out by Matthew Hill in order to determine the time at which the proton source target plasma started interacting with the wire, this figure only shows data from before that time.

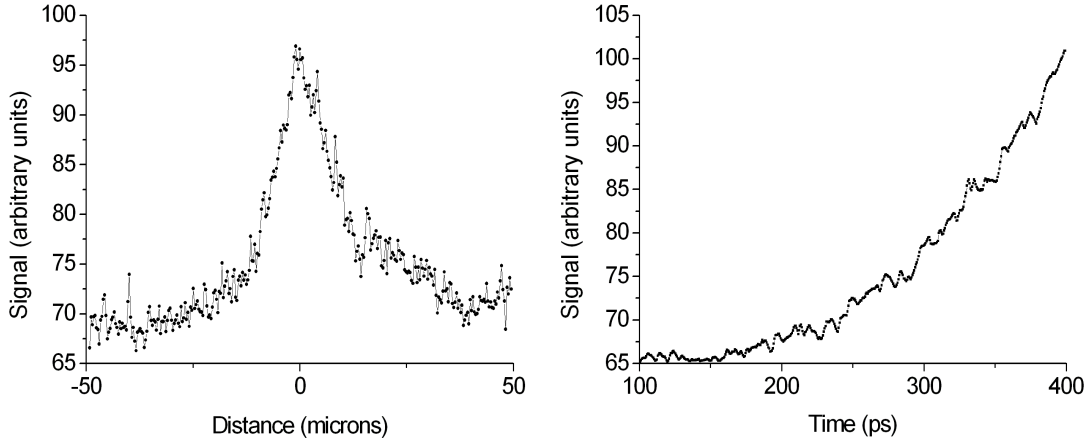


Fig. 4.12: Using Fig. 4.11 the data was integrated and averaged in order to produce spatial and temporal lineouts. a) Spatial lineout across the wire integrated over 300 ps of emission. An approximately symmetric signal is observed implying uniform heating of the wire. A slight increase is observed on the rear surface which may be attributed to background noise. b) Temporal lineout of emission integrated over 50 μm either side of the wire. From the start of emission at 165 ps the emission continues to rise over the timescale of the measurement.

of the imaging system is on the order of the initial wire size so any slight non-uniformity may not appear. However, we can use the approximate symmetry of the spatial lineout as a qualitative indication of near uniform heating between the front and rear surface. The slight increase on the rear surface may be a result of background noise increase from scattered light - with only one shot and a low signal to noise ratio further more detailed analysis is not warranted.

The temporal lineout shows an unexpected increase in emission over a long time period (~ 300 ps). This does not necessarily correlate with an increase in ‘real’ temperature as the expanding wire may only be offering an increased detectable volume of emitting gold at a near constant temperature as it begins to expand. However, we would expect any such change to be gradual. On closer analysis Fig. 4.12 appears to show approximately 100 ps of linear increase and then an acceleration beyond the linear case indicating a two-stage process. While the first 100 ps may be a result of electron-ion temperature equilibration the emission would afterwards be expected to increase at a slower rate. The increase in rate could be an indication of a slow, non-isochoric, heating process potentially caused by high Z ions.

4.4 Discussion

In order to explain the discrepancy between the non-uniform energy deposition calculated from our ion spectrum and the near uniform expansion shown in the XUV data it is necessary to analyse the time duration the heating occurs over. Although the ions are emitted on a time period of the order of the driving laser pulse, the spectrum we have measured will give rise to a time of flight dispersion over the distance to the wire target. Fig. 4.13 shows the results of a simple $time = \frac{distance}{speed}$ calculation for the ions we detect over the $250\text{ }\mu\text{m}$ target separation. This separation is an approximation as the accelerating field will be varying and applied over the rear surface plasma which extends for 10’s of microns.

The first key observation from Fig. 4.13 is that the heating is not isochoric. Thus the Foord method is not strictly applicable as the expansion is not isentropic. The fast protons, which do not completely stop and thus do not have a Bragg peak in the sample (>0.8 MeV for $10\text{ }\mu\text{m}$ of aluminium), and as such deposit energy at an approximately constant rate throughout the sample, all arrive within the first 20 ps. After this time the slower protons Bragg peak towards the front of the sample with the slowest arriving

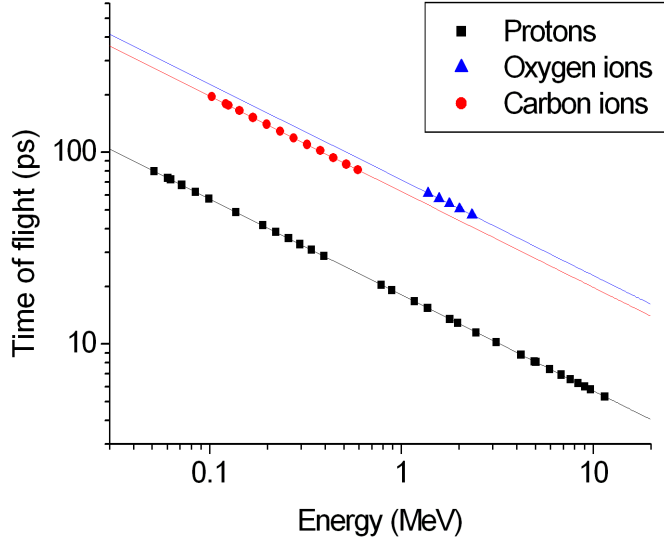


Fig. 4.13: Graph showing time of flight dispersion as ions detected by the magnetic spectrometer propagate over the source to sample separation distance (250 μm). While the ions are emitted over a period of ~ 1 ps, by the time they reach the sample there is over 200 ps of effective heating. In particular the low energy protons and vast majority of carbon and oxygen ions which deposit all their energy in the front surface of the sample take 100-200 ps to arrive.

at ~ 80 ps with any significant fluence. Fig. 4.7 indicated the energy deposition is still approximately uniform at this point, though slightly skewed towards the front surface of the wire. However, for the following 200 ps the heavier carbon and oxygen ions begin to arrive and these do not penetrate beyond ~ 2 μm at solid density.

At this point it is necessary to change our analysis in two important ways. We need to change from a one dimensional sample analysis to a two dimensional one and also take into account the density change as the later ions arrive and the wire has expanded.

When the first carbon and oxygen ions initially arrive at ~ 100 ps Fig. 4.2 shows that the wire has expanded to at least twice the original size with, as expected, the lowest densities on the sample surface. These heavy ions then deposit all their energy within the first ~ 3 μm of expanded aluminium - here they travel twice as far as expected for cold, solid aluminium in the reduced density. At this point, we have a situation where our expanding wire has a few micron hemispherical shell, facing the primary target, which experiences a significantly higher heat flux than the core for the time period of 100-200 ps after the onset of heating. In order to gain an insight into the order of magnitude time scale for heat transfer to the core we solve a simplified one dimensional heat equation with several assumptions.

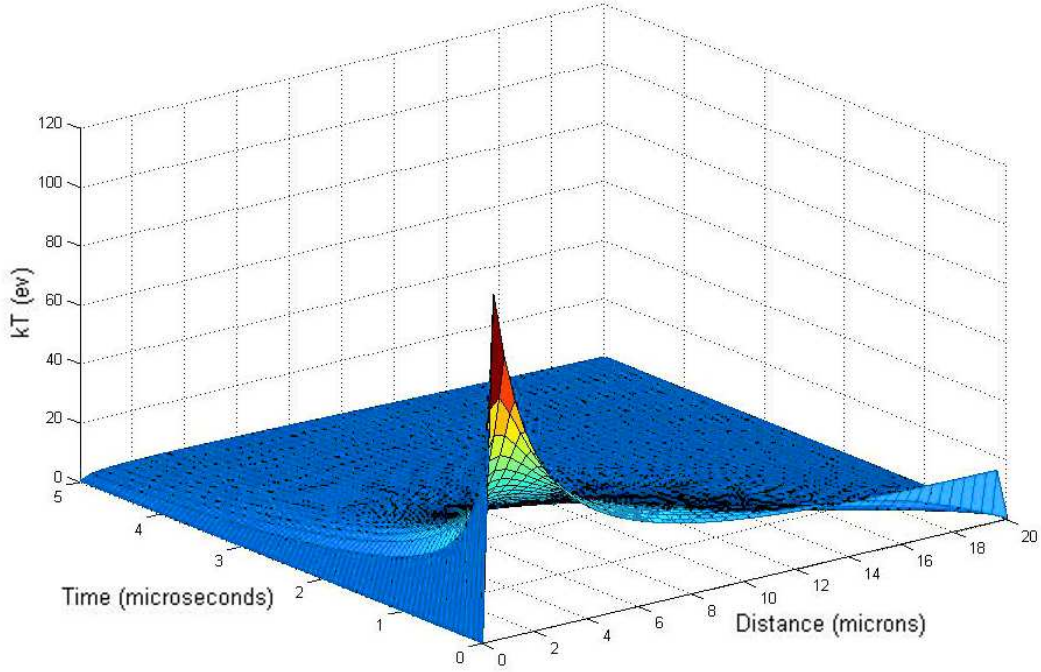


Fig. 4.14: Graphical representation of the numerical solution of the 1D heat transport equation. It is clearly seen that heat diffusion occurs over hundreds of nanoseconds and will not transport heat from the heavy ion heated shell to the core on the ~ 100 ps time scales we investigate. The solution also suggests that the shallow temperature gradients will not induce a shock.

The one dimensional heat equation is:

$$\frac{\partial T}{\partial t} = \frac{k}{c_p \rho} \left(\frac{\partial^2 T}{\partial x^2} \right) \quad (4.10)$$

where T is the temperature at time t along the axis x for a sample with specific heat capacity c_p , density ρ and thermal conductivity k . For simplicity we use the specific heat capacity and thermal conductivity of aluminium at room temperature - $0.897 \text{ Jg}^{-1}\text{K}^{-1}$ and $2.35 \text{ Wcm}^{-1}\text{K}^{-1}$ respectively. In a plasma the conductivity is significantly higher so our time scale will be an overestimate. Again to simplify the variables in the calculation a linear density ramp is assumed rising from zero to solid density at the core, 2.7 gcm^{-3} . Starting with an initial temperature distribution $T(x,0) = f(x)$ using the calculated energy distribution in Fig. 4.8 with $\Delta T = \Delta Q / c_p \rho$ as seen in Fig. 4.14 at $t = 0$. Dirichlet boundary conditions (which specify the value of the solution at the boundaries) of room temperature on the sample edges are also set. The results of this calculation are shown in Fig. 4.14.

Although these are very much limiting case calculations and approach the problem

from a condensed matter viewpoint it gives us an order of magnitude estimate of the time it takes for the high temperature shell to heat the core - several hundreds of nanoseconds. This is a static density case. In reality the sample will be expected to expand more, so that the reduced density will dramatically speed up heat transfer. However, all these mechanisms occur on a much longer time-scale than we investigate.

An analysis of heat conduction from the other limiting case of the plasma viewpoint, which is more applicable at early time where the $k_B T$ is well above the Fermi energy, was carried out by AWE using 2D cylindrical NYM rad-hydro code calculations with electron and radiation heat transport. Here the Spitzer heat conduction is non-linear and depends on $T^{5/2}$ and ∇T [21]. This analysis also shows no significant heat transfer on our time scales [113]. The energy transfer is solely driven by the pressure gradients and hydrodynamic motion. As such, the high energy shell left by the non-uniform ion heating as seen in Fig. 4.15 will stay proportionally localised as the sample expands. This means that if the core, which drives the expansion, and the rear half of the wire are quasi-uniform we can ignore the front surface as it will have little interaction with the rear surface over hundreds of picoseconds.

The author suggests two possible explanations for the difference between the near uniformity seen in the XUV data and the non-uniformity predicted by the deposited energy calculations. Firstly, the slow ions arrive after significant expansion has already occurred, driven by the pressure from the core. Thus the expansion has something akin to ‘inertia’ which is only lightly affected by the late ions whose energy remains at the low density surface, and does not exert much pressure. In order to prove this hypothesis more detailed temperature measurements are required. Although our UV temperature measurements suggest an increase in temperature they do not have the resolution to show conclusive results. Further experiments plan to improve on this. Secondly, we have so far assumed a classical time of flight from the source to the sample. However, very little is understood about the process occurring within the tens of microns over which the ions are accelerated from the surface plasma on our gold disc. We know that it is not a point source but the whole 200 μm disc acts as a source. Literature also suggests that electron, and to a lesser extent ion, refluxing occurs [120], which may act to provide a turbulent acceleration process from a large volume as the degree of refluxing varies over the duration of the laser pulse. Thus the spectrum we measure at our detectors may be significantly different from that which heats the sample at 250 μm from the foil.

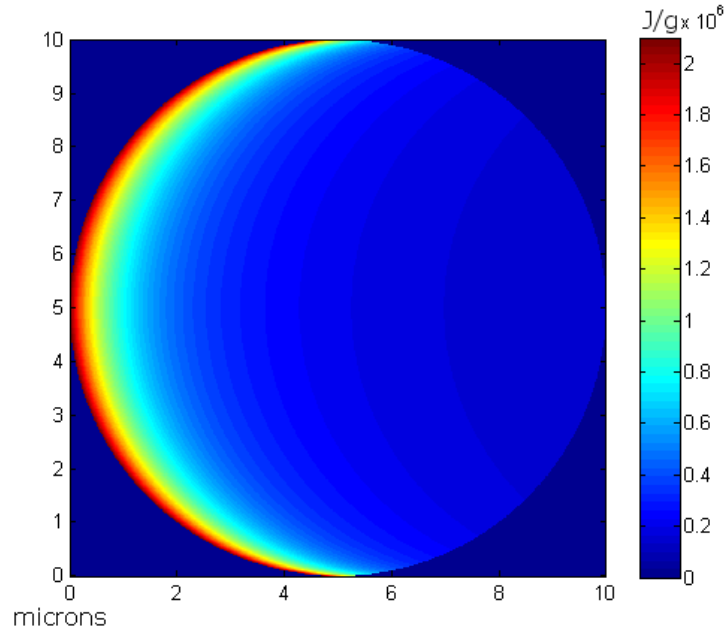


Fig. 4.15: 2D plot of calculated total deposited energy in Al where the ion source is planar and to the left of the figure calculated by the author using SRIM. The wire is shown unexpanded. In real terms much of the front surface deposition would arrive after the wire has expanded to more than twice the original size. This is not shown as the values for stopping power in this WDM regime are unknown and the expected relatively slow heat transport suggests that the deposited energy will remain proportionally localised on our time scales.

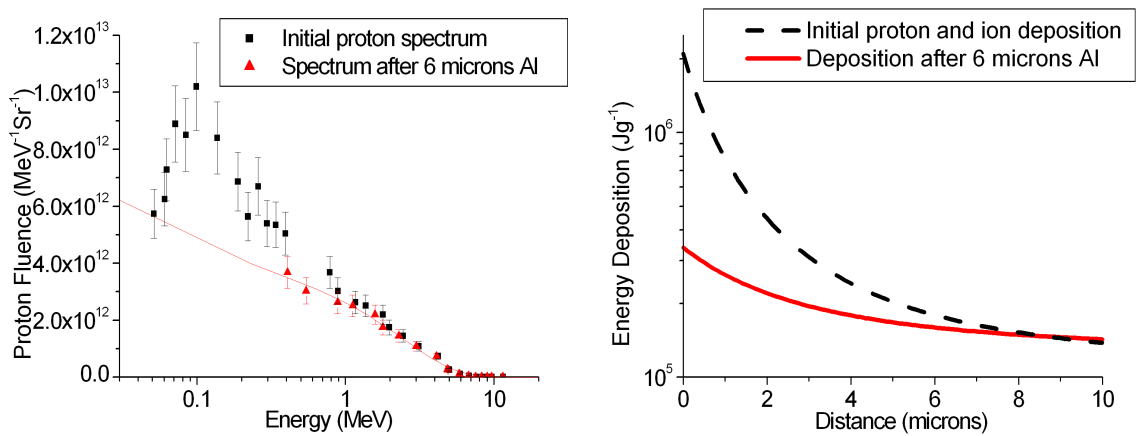


Fig. 4.16: Ion spectrum and energy deposition with, and without a 6 μm aluminium ion filter. Energy loss was calculated using cold stopping powers from SRIM and applying these to the unfiltered beam measured in shot 56. The ion filter is expected to stop all carbon and oxygen ions we detect and all protons below 0.5 MeV.

In order to increase uniformity of heating on one shot we attempted to use a 6 μm aluminium filter to block the heavy ions and the slow protons (< 500 keV). Unfortunately, no spectrum diagnostics were accessible on this shot. However, we can use SRIM to predict what the altered energy deposition would look like, as shown in Fig. 4.16.

To compare this calculated deposited energy with the XUV data shown in Fig. 4.5 we can consider the ratios of front to rear surface expansion velocity to the deposited energy ratio. Here, we use the adiabatic sound speed, $C_{sa} = \sqrt{\frac{\gamma P}{\rho}}$ and at the boundary of the isentropic expansion the speed is $v_b = 2C_{sa}/(\gamma - 1)$. Hence the squared ratio of the front and rear boundary speeds estimates the ratio of the front and rear pressures, which is approximately proportional to deposited energy. The XUV data in Fig. 4.5 shows a squared ratio of 3.5 whereas the deposited energy has a ratio of 2.5. This suggests the ion filter may be thinner than stated or contain pinholes. Another, more attractive possibility, is that the filter, which is equidistant between the ion source and sample, acts to focus the electric field generated by the disc target. Thus the outside of the ion beam, which has a higher proportion of low energy protons and does not typically see the sample now provides additional heating and increases the non-uniformity. This can be thought of as an effect similar to the method of images solution for a charge above a conducting plane - essentially using the fact that the field component parallel to our conducting foil is zero. Quantitative analysis of this is not trivial due to our lack of detailed knowledge of the initial acceleration process and the shape and time varying extent of the electric fields present.

The energy deposited is consistently lower than that calculated from the XUV data - typically by half as much. Much of this may be a result of systematic error yet to be discovered - such as the limited dynamic range of the ion spectrometer leading to an underestimate of total beam energy. However, the author believes this difference may for a large part be a result of the unknowns of warm dense matter physics - the purpose of our investigation. Numerical models such as NYM are based on plasma parameters and have increased error in the density space we investigate. Also, throughout these results and discussion we have used cold, solid density stopping powers from SRIM. Clearly as the ions heat and ionise the sample, the electronic and nuclear stopping will be significantly altered and thus the deposited energy may change significantly. Further experiments are vital to obtain data on this little known subject. This is the topic of chapter Ch. 6.

4.5 Summary

The main topic of this chapter was the experimental development of the generation and characterisation of warm dense matter conducted at AWE and the Rutherford Appleton Laboratory. The data presented included temporally and spatially resolved XUV density measurements of a proton heated wire expansion simultaneously with its UV emission. These were analysed with a view to comparing them with the measured ion spectrum. This comparison of deposited energy calculations disagreed by a factor of two, which highlights the need for an experiment to determine the stopping power of warm dense matter. This is expected to be substantially different to the cold stopping powers used for this analysis. The ion spectrum measurement was designed to accurately measure low energy protons below the detection threshold of a radiochromic film stack (<4 MeV) and heavier carbon and oxygen ions. These particles are believed to lead to non-uniformity of heating which impacts on equation of state measurements.

With a view to evolving the method with future experiments on the new Orion laser, improvements on the UV emission imaging diagnostic would move to a reflection based focusing system to increase resolution and eliminate chromatic aberration. Accurate calibrated pyrometry could then be used to determine temperature. Lower magnification will give a better signal to noise ratio. Also, to accurately determine deposited energy over the time scale of the 200 ps ion time of flight dispersion, work needs to be carried out to investigate the stopping power of warm dense matter in order to accurately benchmark models. Finally, the use of an ion filter in future experiments was discussed with the view to increasing the uniformity of heating by eliminating ‘cold’ or heavy ions which would otherwise preferentially heat the front surface of the target.

Chapter 5

High Energy Density Blast Waves in Cluster Gases

This chapter describes laser driven shock experiments based on the theory and experimental methods described in chapters 2 and 3. The experiments described here were conducted in 2007 on the Vulcan laser (Target Area West) at the Central Laser Facility in the Rutherford Appleton Laboratories and throughout 2007-2011 on the Imperial College Laser Consortium's Glass laser system. The chapter will begin with a discussion of two drive beam geometries used to investigate the upstream parameters of the expanding shock. We then go on to detail the progress towards proton probing of these shocks in order to resolve transient electromagnetic fields. Finally, high energy scaling of single blast waves on the Vulcan laser is described for gases of varying atomic number with particular reference to the continuing high absorption and unique spatial structures at these energies, in conjunction with a large radiative precursor.

5.1 Probing Upstream Material via Two Beam Geometries

The following section details experiments on the Glass Laser carried out with the aid of S. I. Olsson Robbie. The high laser energy absorption properties of clusters ($>90\%$) compared to typical gases ($<5\%$) is used to launch a strong shock, and then probe the material ahead of the expanding shock wave with an additional, perpendicular heating beam. The motivation for this arises from astrophysicists querying the applicability of scaling laboratory based experiments to their domain if the ambient material consists of clusters rather than a monatomic gas. One potential effect of this would be a change in the dynamics of a radiative blast wave expanding into a cluster gas, as their increased absorption, as opposed to typical gas, may increase the precursor temperature. This in turn will lead to altered cooling dynamics as the smaller precursor reaches a higher temperature and more energy is reabsorbed to produce a faster shock.

A useful secondary, spin-off experiment with this geometry allows us to study the effect of pre-heating the upstream material on shock dynamics.

This work is currently submitted for publication to the journal *High Energy Density Physics*.

5.1.1 Previous Work

Prior to carrying out this experiment previous work such as that by A. S. Moore *et al.* [102] which describes the radiation emitted by the propagating shock front led us to intuitively suspect that this radiation disassembled the clusters ahead of the blast wave. This theorised mechanism would have an impact on the calculated absorption depth of the emitted radiation and how it is reabsorbed back into the shock as clusters may have a different optical depth to that of standard, or even lightly ionised, gas.

To my knowledge only two previous experiments have been carried out using multiple, colliding, laser driven blast wave geometries in cluster media. J. Lazarus uses a Fresnel bi-prism to produce two near parallel blast waves that expand and collide showing density enhancements at the point of collision [27]. These effects are also observed in our experiment. D. Symes makes use of a second heating beam through an interferometer to destroy clusters in an evenly spaced lattice [121]. The original heating beam then produces regular spaced expanding blast waves which collide with the hope of seeding the Rayleigh Taylor instability.

The 2007 Vulcan laser experiment also attempted to heat upstream material with laser

generated ~ 50 eV soft x-rays in order to assess the affect of radiative heating on blast wave velocity. The analysis of this experiment was carried out by M. Hohenberger and is detailed in his thesis [28]. An increase in velocity is demonstrated as the blast wave passes into the preheated material. However the level of heating is difficult to quantify with the limited data set currently available.

5.1.2 Experimental Setup

The Imperial College Laser Consortium Glass laser, detailed in Fig. 3.5, was used to produce pulses of ~ 1 J at 1054 nm with 1 ps duration. Fig. 5.1 illustrates how this output beam was divided to produce two ~ 500 mJ variable delay heating beams and a low energy green probe beam. Producing the beams optically from one initial beam ensures they are temporally well synchronised and do not suffer from electronic timing jitter. The probe beam on reaching the sample contains two polarisation states where one is delayed in time with respect to the other. This allows probing on one line of sight at two different times for a single shot. This is significant as the experiment is relatively low repetition rate (0.01 Hz) with a varying reproducibility due to the varying laser energy.

Fig. 5.2 is a simplified representation of the three orthogonal beams hitting the cluster target produced by the modified gas jet (Fig. 3.23). The two heating beams are focused with $f/3$ GRIN (gradient-index) lenses which have a focal length of ~ 13 cm. The GRIN lens uses an axially varying refractive index to eliminate the spherical aberration occurring in typical lenses. The green probe image rays are collected and collimated with a large $f/4$ lens.

The three beams were spatially overlapped at low power using obscuration with a $10\ \mu\text{m}$ wire, and temporally overlapped to within a picosecond using the vertical heating beam as a reference for t_0 . Whereas traditionally timing is carried out with a fast streak camera here a different method is used. The 0.5 ps duration probe beam is timed to the vertical beam by moving the optical delay stage to bracket the timing such that the fast initial ionisation stage of the cluster gas is only just captured by the imaging cameras. The horizontal beam is then timed to the vertical beam using the cluster break-up effect we are investigating. If the horizontal beam arrives later than the vertical beam it's resulting blast wave will have a gap in the middle where clusters are destroyed before its arrival, thus lowering the absorption. Similarly if the vertical beam arrives later *it* will have a gap in the middle. In this manner the beams can be timed such that they appear to merge

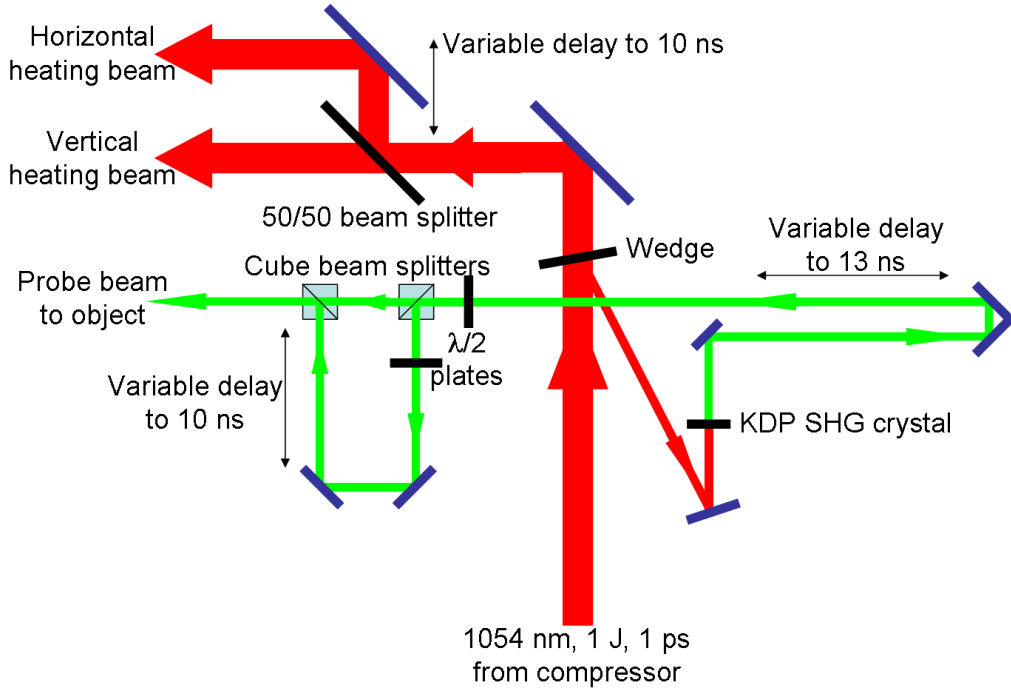


Fig. 5.1: Schematic of beam splitting to provide dual heating beams and a time delayed dual polarisation combined probe beam. A wedge is used to select 4% of the main beam which is then frequency doubled to produce green light that travels through a variable delay slide. This green light is then split 50/50 with a polarising beam cube and one polarisation state is delayed before they are recombined with another cube. After passing through the sample the two polarisations can be split and thus two separate times can be recorded. The main heating beam is divided with a 50/50 beam splitter and the horizontal beam is delayed with respect to the vertical beam which hits the target at t_0 .

into a cross as seen in Fig. 5.5b which remains to ± 1 ps. The reader should note that a fine translation stage is necessary in the variable delay as to delay a beam by 1 ps requires a movement in the delay of 0.3 mm.

The gas nozzle was fired with an opening time of 4 ms and timed such that it had been open for 3 ms before the heating beams arrive. This allows the jet to stabilise as there will be a reduced onset of clustering in the initial turbulent flow. These experiments were carried out in argon as it is a relatively cheap noble gas with a reasonably high atomic number ($Z=18$) which clusters easily. The gas jet backing pressure and temperature parameters were chosen such that later experiments with different gases could be carried out with a similar average number of atoms per cluster as seen in table 5.1. Ideally the parameters would be chosen to have similar average mass densities however this necessitates impractically high or low pressures and as such a compromise is made. However, the high Hagen parameter, Γ^* , ensures large clusters with high laser absorption. Hydrogen is the possible exception to this rule as some experiments demonstrate lower laser absorption -

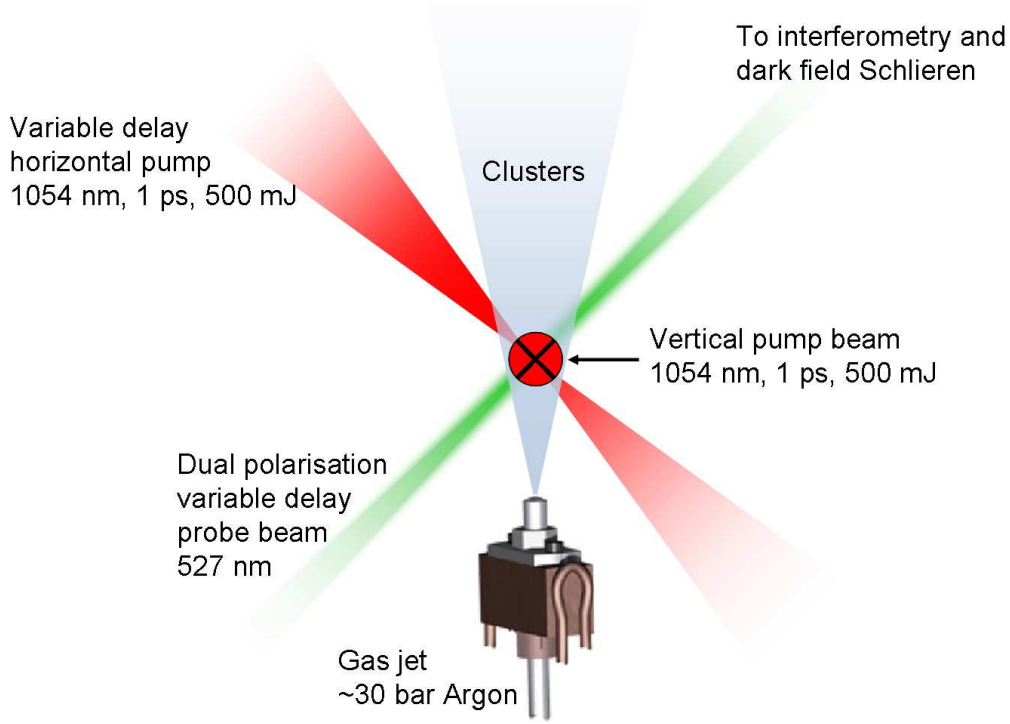


Fig. 5.2: Top view of the experimental setup for perpendicular blast wave experiments. Here we see that three laser beams - two pump and one probe - arrive at the cluster target from three perpendicular directions. These beams need to be overlapped in space and, at least initially, in time to within 1 ps - a non-trivial task. The modified, extended gas jet nozzle allows it to be within ~ 2 mm of the overlap point of the probe and heating beams without blocking any light.

$\sim 50\%$ [28] - though this may be due to the faster cluster disassembly time for light gases.

Gas	$T(K)$	p (bar)	p_{ideal} (bar)	$\Gamma^* (\times 10^3)$	$N_c (\times 10^3)$	R_c (nm)	γ
Ar	298	30	30	21	42	7	1.67
H ₂	123	35.5	245	21	42	6	1.59
Kr	298	17.1	14	21	42	8	1.67
Xe	298	9	9	21	42	9	1.67

Table 5.1: Experimental gas jet parameters for similar conditions in different gases. The pre-expansion nozzle temperature, T , and backing pressure, p , are those required to produce an equal value for the Hagen parameter, Γ^* , and therefore mean average number of atoms per cluster, N_c . The cluster radius is then R_c . The ideal pressure, p_{ideal} , is that which would give an equivalent mass density distribution to an argon reference value but this is impractical with our gas nozzle for H₂. The polytropic index values, γ , are shown for completeness and are taken from [122].

The experiment was carried out by delaying the horizontal heating beam with respect to the vertical beam. The distance the delay stage was moved was converted to a time

and as such had an error of ± 3 ps. Similarly, at each heating beam time delay the probe beam delay was scanned.

After the probe beam passed through the sample and the image rays were collimated by the collection lens they were magnified ($\times 2$) then split via a 50/50 beam splitter so that simultaneous interferometer and dark-field Schlieren images were taken with the polarisations split at two time delays as shown in Fig. 5.3. Use of a resolution calibration slide showed the spatial resolution was at best $10\text{ }\mu\text{m}$. The Schlieren stop was typically the tip of a $\sim 1\text{ mm}$ wire - or an unravelled paper-clip.

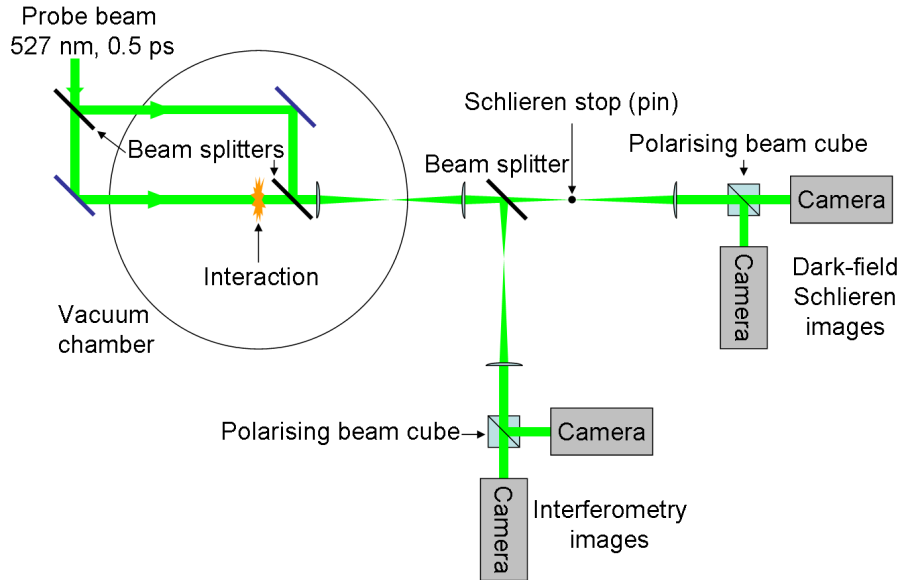


Fig. 5.3: Schematic of imaging setup for colliding blast waves experiment. The beam geometry did not allow use of the Michelson interferometer so a Mach-Zehnder was used. A three lens imaging system was used to magnify the beam ($\times 2$) and provide a focal plane outside of the vacuum chamber to place the Schlieren stop in. A 50/50 non-polarising beam splitter divides the beam into a Schlieren arm and an interferometry arm. A polarising beam splitter in each diagnostic is then used to divide the two time delayed polarisations to form a focused image on individual Sumix M72 CMOS USB cameras which are connected to a computer.

Multiple shots were taken at each delay to obtain an average. The fluctuations are accounted for mostly by the shot to shot energy variation in the regenerative amplifier of the Glass laser.

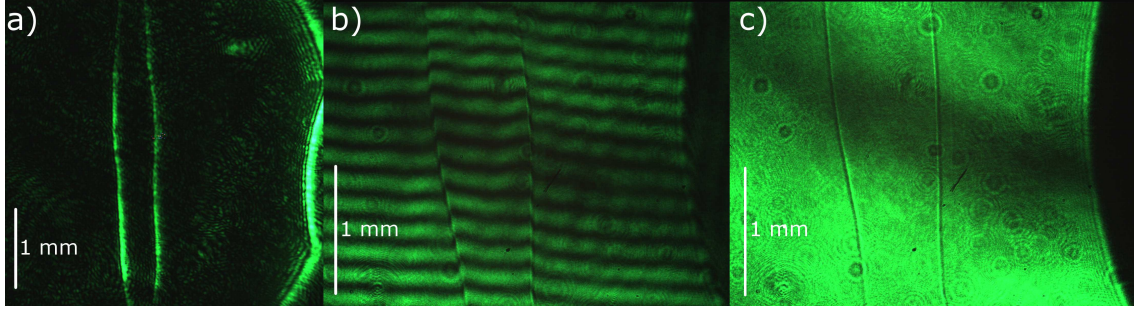


Fig. 5.4: Raw dark-field Schlieren (a), interferometry (b) and shadowgraphy (c) images of a vertical blastwave. The green probe arrives 10 ± 0.1 ns after t_0 . The shadowgraphy image was obtained by blocking one arm of the interferometer. The images are from similar shots with no horizontal heating beam and 30 bar backing pressure of argon. The vertical beam contained ~ 400 mJ of energy. The circular dots in b) and c) are dust on the camera chip and can be ignored as artefacts. The round shadows in the right of each image is the gas jet nozzle and is included for size comparison.

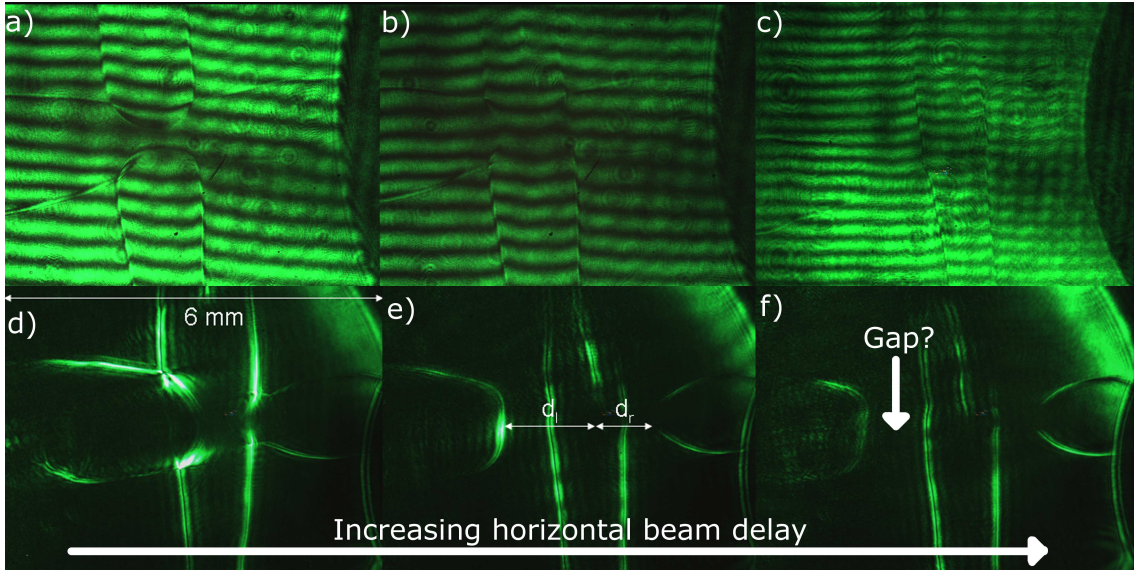


Fig. 5.5: Raw interferometry (a-c) and Schlieren (d-f) images of a horizontal beam delay scan. All images were captured 10 ns after the vertical beam arrives. a) Horizontal beam arrives ~ 30 ps before vertical beam. b) Horizontal and vertical overlap in time. c) Horizontal arrives after 3 ns. The interferometry images are not ideal for this cross beam experiment as it only allows for analysis of one of the beams. For this reason Schlieren and shadowgraphy imaging was very useful to resolve the blast wave edges. d) Horizontal arrives 10 ps after vertical. Bright spots indicate density enhancements as the two beam collide. e) Horizontal arrives after 2.3 ns and f) after 4.2 ns. d_l and d_r indicate the measured separation from the centre of the vertical blast wave to the edge of the horizontal. d_r is smaller than d_l as the density is higher closer to the gas nozzle (shown in shadow to the right) and exerts a higher stopping power.

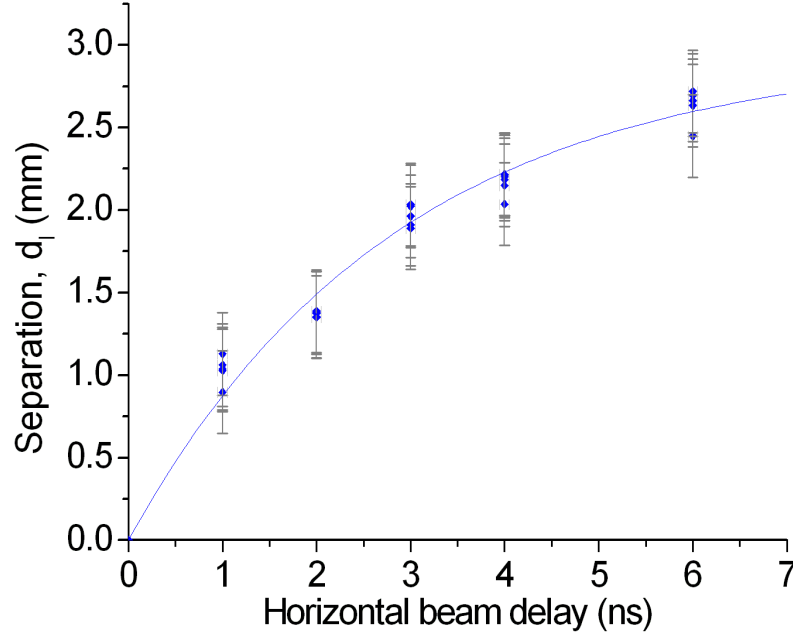


Fig. 5.6: Graph of horizontal beam delay against left separation. Left separation, d_l , is defined as the distance from the vertical beam centre to the edge of the left blast wave. The two main sources of error arise from shot to shot variation in laser energy (though blast wave velocity is a relatively weak function of deposited energy, $\sim E^{1/4}$) and difficulty in accurately locating the edge of the left blast wave on some shots.

5.1.3 Results

A complete time scan of argon was carried out. Fig. 5.4 shows the types of raw dark-field Schlieren, interferometry and shadowgraphy images obtained with just a vertical blast wave. Unfortunately, time and equipment constraints did not allow full scans of other gases except for a limited number of shots in hydrogen. Fig. 5.5 shows interferometry and Schlieren images as the horizontal beam is delayed. As the horizontal beam is increasingly delayed a large gap where no blast wave exists rapidly grows. We can plot a graph of delay against size of gap where its extent is defined to be from the centre of the vertical blast wave to the edge of the horizontal as seen in Fig. 5.6. We plot the gap to the left of the vertical blast wave, d_l , as it is more pronounced than the one on the right, d_r . We account for this difference due to an increase in density nearer to the gas jet.

As the delay increases a ‘bump’ can be seen to grow on the edge of the vertical blast wave as shown in Fig. 5.7. Looking at the phase shift in Fig. 5.8a we see an increased electron density between these bumps. The author attributes this to the blast wave ex-

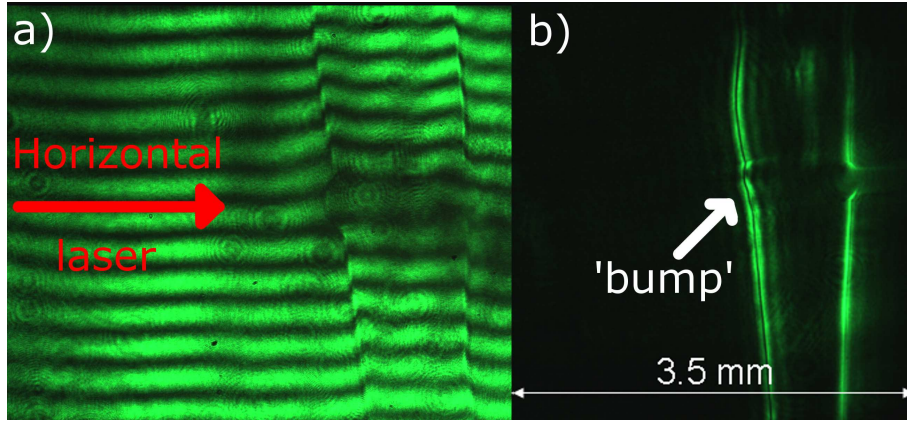


Fig. 5.7: Interferometry and Schlieren images of a blast wave moving into pre-heated material. Here the horizontal beam arrives 7.9 ns after the vertical and is off the screen. A bump can be seen in the blast wave where the laser energy has pre-heated the gas. The bump is more pronounced as horizontal beam delay increases suggesting a faster blast wave velocity into the pre-heated material.

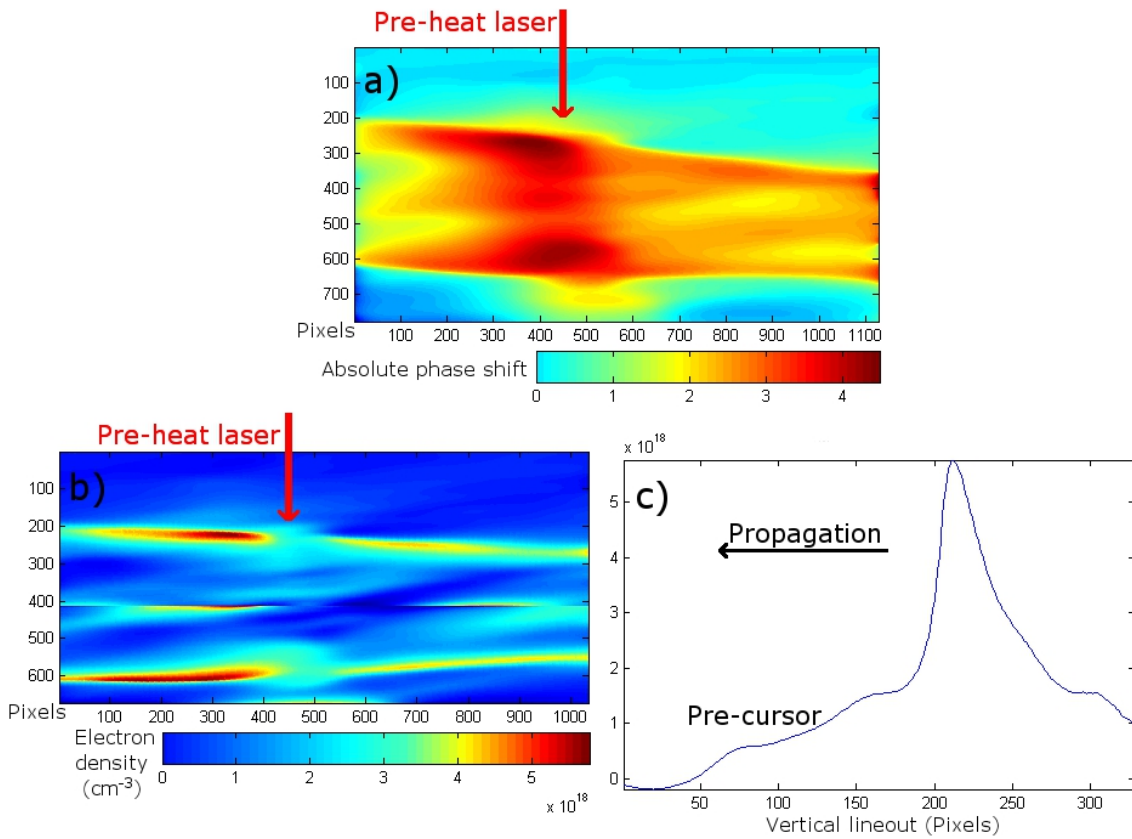


Fig. 5.8: Analysis of pre-heated interferometry image Fig. 5.7 carried out using the PAT software. a) Phase image of the vertical blast wave where the increased density can be seen as the horizontal laser passes through. b) Abel inversion of the phase image. This is not properly applicable in the region of the horizontal laser as the cylindrical symmetry is broken. The line through the centre of the blast wave is an artefact arising from the inversion process. c) Vertical lineout taken from the right edge of the Abel inversion showing the top half of the blast wave. A small radiative pre-cursor can be seen to the left of the peak that extends for ~ 0.5 mm into the upstream material.

panding into pre-heated material around this area. Although, no horizontal blast wave is observed the laser is still present and deposits energy. This energy deposition must be much reduced otherwise a blast wave would form. To quantify the energy deposited vertical interference fringes are required in order to recover the ionisation state of the gas and hence an estimate of the energy needed to achieve this. This lower early energy deposition also accounts for the larger ‘bump’ at larger horizontal delay as less laser energy has been lost, which would have been absorbed to form a blast wave, by the time it hits the vertical blast wave.

Analysing the vertical beam interferograms away from any horizontal beam interference as shown in Fig. 5.8 we see a radiation pre-cursor does exist to the left of the peak but only extends for approximately 0.5 mm - much less than the observed d_l .

5.1.4 Discussion

The unexpected result shown in the previous section is that there is a cluster disassociation ‘wave’ moving well ahead of the radiative pre-cursor, which destroys the upstream clusters, as shown by the lack of the perpendicular blast wave in this region. Fig. 5.5a shows that, unlike the radiative pre-cursor seen in Fig. 5.8c, the source of this disassociation wave is not the shock front but lies behind it. This suggests the disassociation wave is sourced from the initial ionisation event when the laser first forms a plasma filament in the highly absorbing clusters.

In this initial ionisation event three forms of radiation can cause a moving disassociation front - photons, electrons and ions. Photons travelling at the speed of light can be dismissed as our observed disassociation wave travels significantly slower than this. This leaves electrons and ions as possible sources of the disassociation wave. Fig. 5.9 plots the velocity of the disassociation front and the possible energies of the electrons and ions from this velocity.

In order to predict the maximum energies of electrons and ions produced in the initial cluster explosions we return to the nanoplasma model. The use of the nanoplasma model is valid here as we have relatively ‘long’ pulses and heavy atoms. For ultra short pulses (<100 fs) and low Z atoms the model has been shown to have significant error [123]. The laser intensity incident per cluster is $\sim 10^{17}$ Wcm $^{-2}$. The model predicts maximum electron energies of hundreds of keV which is verified by experimental measurements [124] [125]. This is much higher than the 3 eV energies shown in Fig. 5.9 and as such we

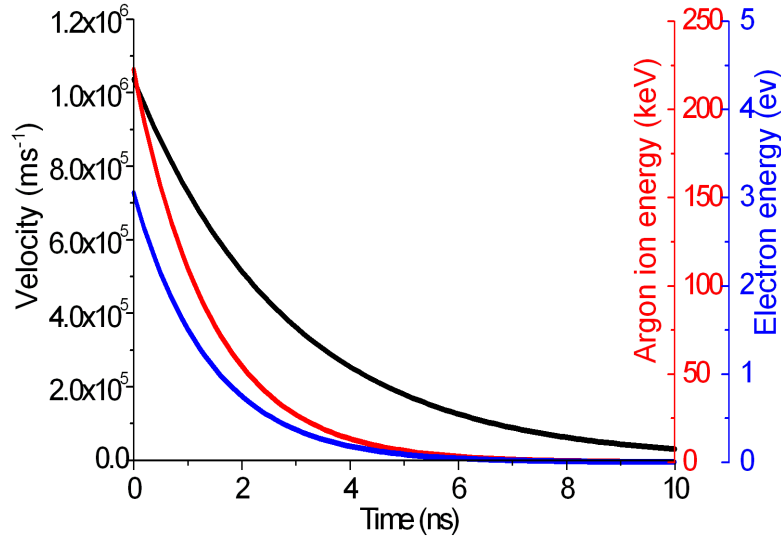


Fig. 5.9: Graph of disassociation front velocity against time. Here the gradient of the best-fit curve from Fig. 5.6 has been used to obtain velocity (black, upper line). As a guide to determining the origin of the disassociation front the velocity has been converted to energy for argon ions (red, middle curve) and electrons (blue lower curve). This shows an initial ion energy of 222 keV and an initial electron energy of 3 keV.

discount electrons. However, as shown in Fig. 5.10 and in experimental measurements [126] [127], ion energies of ~ 100 keV are also predicted. This suggests that ballistic ions launched in the initial laser-cluster interaction are the source of the disassociation wave front. The paper by Kagawa *et al.* clearly demonstrates that argon ions with energies as low as 3 keV will disassemble clusters [128]. Kagawa also suggests that after the 100 ps cluster destruction the atoms are not left ionised as a warm plasma but exist as a neutral, monatomic gas.

Having established that ballistic ions are the cause of the cluster disassembly ahead of the blast wave it is necessary to account for the decrease in velocity of this disassembly front. The stopping power of these upstream argon clusters will cause a multi-keV argon ion to lose energy as it propagates. Although the stopping power of clusters has not to my knowledge been explicitly examined previously we can attempt to model it as a typical gas [129]. The SRIM model using ZBL stopping (Ziegler, Biersack and Littmark [130]) which includes electronic and nuclear stopping can model this. At these relatively low ion energies nuclear stopping can become more significant than electronic stopping as the chance of a core collision increases. The initial ion energy of ~ 222 keV estimated from the observed disassociation front velocity at early times is used and propagated through an

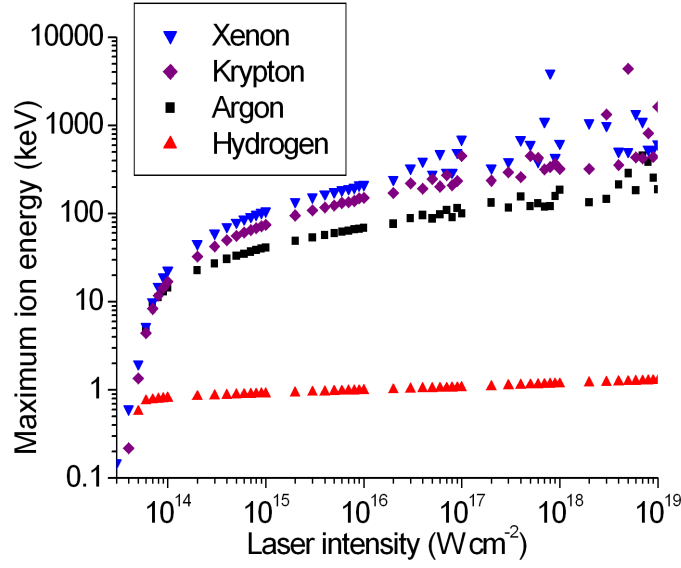


Fig. 5.10: Nanoplasma model result for maximum ion energy from argon, hydrogen, xenon and krypton clusters at our gas jet conditions. The laser intensity incident per cluster for our experiments is on the order of $10^{17} \text{ W cm}^{-2}$. It should be noted that the nanoplasma model is unreliable at modelling hydrogen. The clusters sizes are chosen to have the same number of atoms per cluster, N_c .

argon gas of decreasing density as described by the hydrodynamic gas expansion model (Eq. (3.19)) with an initial distance of $\sim 2 \text{ mm}$ below the nozzle. Fig. 5.11 shows the results of this theoretical fit.

It should be noted that this model does not differentiate between initial ion charge states. This comes about for two reasons. Firstly, Ziegler tells us that it is the *effective* charge state which is a function of velocity, gas pressure and electron screening which is used in the code [131]. Secondly, the nominal charge state of the ion fluctuates dramatically as it transits the material and interacts via electronic and nuclear stopping processes. Indeed, often the charge state is higher on exiting the material [132].

Fig. 5.11 shows that SRIM models the overall energy trend well. However, SRIM underestimates the stopping power of argon clusters. This may suggest that clusters have a higher stopping power than an un-clustered gas. To accurately model the local solid density but average gas density properties of clusters we need to consider the physics behind the stopping power of solids as opposed to gases, and some of the inherent error in SRIM modelling.

Generally, for an identical atom density, gases exhibit a higher stopping power than solids (with the exclusion of effects such as *channelling*). The explanation for this counter-

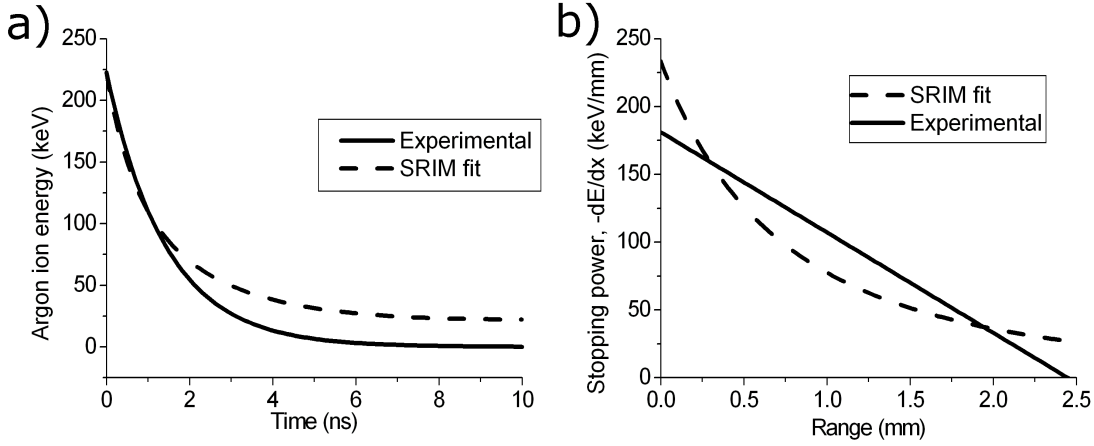


Fig. 5.11: SRIM model to fit the stopping power of argon gas. The SRIM model used densities from the free gas jet expansion model starting at ~ 2 mm below the gas jet. An initial argon ion energy of 222 keV was used. Comparing the energy (a) and the stopping power (b) it is interesting to not that the SRIM fit underestimates the stopping power of the clusters.

intuitive effect is that the ion's polarisation of the solid's electrons act to shield the ion. This lowers the effective charge and reduces distant collisions [131]. However, for an ion passing through a single cluster the solid density of the cluster only exists for a few nanometres before the ion exits the cluster presumably leaving this shielding effect behind. This means that the two, typically less significant, effects in solids which increase stopping power may become dominant. Interactions between the ion and subsequent atoms occur before the ion's electrons completely relax, increasing the effective charge. Also the solid's electrons may form a continuous, mobile cloud which interacts with the ion collectively. In a gas these effects contribute far less to the stopping power but vary with gas density - over a factor of two differences in stopping power have been observed as the pressure moves away from standard pressure (STP) [131].

Ziegler goes on to comment that below 200 keV/amu, the region we investigate, the difference between solid and gas stopping increases by as much as a factor of two.

As the ballistic ion wave sweeps upstream of the blast wave it will disassemble clusters leaving cold argon gas in its wake. As our second horizontal heating beam propagates through this gas it will deposit energy, although significantly less than in clusters of comparable average density. A literature search suggests that due to plasma defocusing we will have intensities on the order of $\sim 10^{17}$ Wcm $^{-2}$ ionising the gas [133]. Literature also suggests that we will ionise the argon gas to a maximum of Ar $^{7+}$ [124]. Thus the vertical blast wave will be expanding into a region of pre-heated gas from which it will absorb

energy and accelerate producing the ‘bump’ shown in Fig. 5.7.

Using this ionisation state of $Z^* = 7$ and an average gas number density of $3 \times 10^{18} \text{ cm}^{-3}$ then ionised gas will have a free electron density of $2 \times 10^{19} \text{ cm}^{-3}$. The interferometry analysis in Fig. 5.8 shows an increased electron density in the region of the horizontal beam. Here, we see a decrease in the centre of the blast wave suggesting that as the horizontal beam arrives 7.9 ns after the vertical blast wave the expansion has swept up mass from the central region, decreasing the number of atoms available for further ionisation and hence the observed electron density.

In order to obtain a more quantitative analysis of the extent of electron densities present and indeed the ionisation state it would be desirable to repeat the experiment with interferometric fringes running in the vertical direction. This would allow an Abel transform to be carried out. Unfortunately time constraints prevented this.

As a proof of principle experiment it can clearly be seen however that pre-heating does allow us to more accurately model an astrophysical scenario. Particularly to model the altered dynamics as a supernova remnant expands into a lightly ionised background gas.

5.1.5 Summary

In summary, this data has shown that the radiative cluster blast wave’s radiation does not, as previously thought, destroy the clusters ahead of the expanding blast wave. Instead a ballistic wave consisting of multi-keV ions produced in the initial laser-cluster ionisation process moves out well ahead of the blast wave with an initial velocity here of 1000 km/s. This ion wave slows down as it interacts with the clusters and disassembles them. Thus the radiative blast wave expands into a uniform, cold or very weakly ionised gas and all the relevant thermodynamic properties associated with it.

The nanoplasma model supports the observed ion velocities of 222 keV from the initial hydrodynamic cluster explosion. Using SRIM to model the stopping power of clusters supports the experimental stopping power trend but suggests that clusters have a higher stopping power than that of typical gas. Further experiments would ideally use an accelerator source to confirm this. However, this perpendicular beam geometry does supply a method of stopping power measurement without the need for an accelerator. An ion spectrometer would verify the ion energies. Further work will also verify stopping powers in different gases of varying atomic number.

This work emphasises the need to tailor the cluster medium with a skimmer to reduce

ion energy loss when using clusters as an ion source.

A proof of principle experiment has also been shown to pre-heat the background gas such that the blast wave expands into a lightly ionised medium as is the case in the interstellar medium. Further experiments would investigate the altered, accelerated dynamics in this regime.

5.2 Proton Probing of a Cluster Gas Blast Wave

To fully characterise our blast wave regime we have established above that the upstream material is a cold, monatomic gas and not clusters. The other question posed by astrophysicists looking to make an accurate comparison to their field is the electromagnetic properties of the blast wave. In particular, the existence of large magnetic fields in supernova remnants has been evident since the detection of their synchrotron radiation [134] [135]. It is also well known that magnetic fields can significantly alter the dynamics of systems such as propagating shocks in astrophysical scenarios.

Indeed, we can speculate that since wherever charge separation occurs there must be electromagnetic fields we should be able to detect these in our blast waves. Of course the strength and duration of these fields will make detection non-trivial. Previous work with laser foil interactions produced collisionless shocks has detected these electric fields [136].

In Ch. 3 we described using protons as a means to probe electromagnetic fields. If we use low energy protons (~ 100 keV) they will detect low strength fields and their short duration (~ 1 ps) allows detection of very transient fields with a high spatial resolution.

Using a cross beam geometry similar to that used in the previous section the horizontal beam was focused onto a $12.5 \mu\text{m}$ aluminium foil and the vertical beam used to create a blast wave as shown in Fig. 5.12. Protons were produced via the target normal sheath acceleration mechanism with an energy of ~ 100 keV (see Ch. 6 for spectrum). The proton beam propagated through a copper mesh (40 wires per inch) to imprint a pattern such that any distortion of this pattern on the CR-39 detector could be attributed to the Lorentz force. The blast wave extent was verified using a green optical probe as described in the previous section. The gas jet backing pressure was increased to 50 bar to increase the size of clusters and blast wave in order to potentially increase the size of any self generated electromagnetic fields. The aim was to scan the variable delay in the proton generating arm to detect varying magnetic fields as the shock evolves.

Moving the CR-39 closer to the blast wave increased the fluence and decreased the

magnification however any divergence would become more pronounced as the CR-39 moves further away. Here, the distance was varied from 20 - 100 mm. Similarly moving the mesh and the proton source in relation to the blast wave will change the volume of blast wave probed and the spatial resolution of the mesh. Typically the source was 18 mm from the blast wave and the mesh 10 mm. The best focus of the horizontal laser on to the foil was established using a high magnification CW imaging system and translating the focusing lens or foil as necessary.

Although we have developed the experimental method we have yet to image a blast wave. Fig. 5.13 shows the results obtained and highlights the problem to be overcome. Fig. 5.13a shows that with no gas, even CR-39 placed at 100 mm still shows a clear image of the mesh. However, with gas added this clear image dissipates even at 25 mm where the divergences have less time to propagate. While at first this seems clear evidence of strong electromagnetic fields unfortunately when the vertical beam is not used, and hence no blast wave exists, the dispersion still remains. As such this is attributed to the mean free path of the protons in the cluster gas being significantly less than the propagation distance through the jet.

Fig. 5.14 shows the results from SRIM of modelling the protons' path through 5 mm of argon gas at a number density of $3 \times 10^{18} \text{ cm}^{-3}$. The increased divergence as a result of small angle collisions can clearly be seen. To overcome this the gas jet backing pressure was reduced to lower the density in the gas jet. This qualitatively began to improve the 'blurring' however not enough to recreate the mesh pattern throughout the jet before the pressure was too low for a blast wave to form.

For future work three options exist to successfully image the electromagnetic fields. Firstly, tailoring the gas jet to reduce the volume of gas the proton travels through and hence reduce divergence. Secondly, use a different, lower atomic number, gas which as seen in Fig. 5.14 for hydrogen will dramatically decrease divergence. Equipment and time constraints prevented us attempting this option. Thirdly, the preferred method, is to move to a higher power laser such as the Vulcan laser. Here higher energy protons (60 MeV), which will exhibit less divergence, can be created and in combination with a RCF stack will image different strength fields (and densities) at different layers. Also a higher energy laser will produce stronger blast waves with potentially larger electromagnetic fields.

Introducing a static magnetic field, for example with a pulsed power solenoid, which

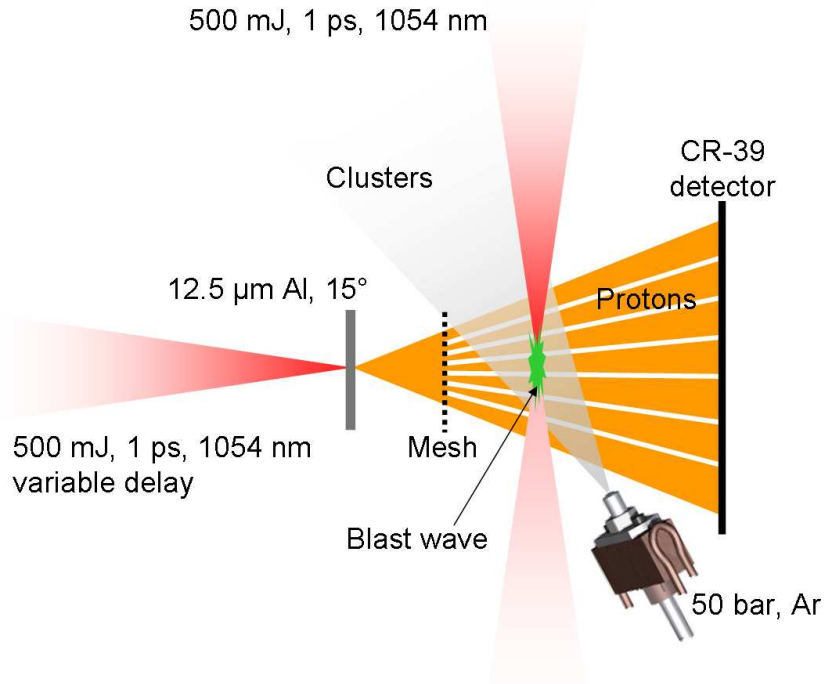


Fig. 5.12: Blast wave proton probing schematic. The main heating beam is split such that one arm produces protons via the target normal sheath acceleration mechanism on a 12.5 μm aluminium foil tilted off normal by 15° . The other arm is used to produce a blast wave in the cluster medium. Not shown here is the use of a perpendicular optical probe to verify the existence of a blast wave. The proton beam passed through a 40 wires per inch copper mesh to simplify our understanding of any deviation of the individual beamlets by the Lorentz force caused by the electromagnetic fields in the blast wave. CR-39 detector was used as it is sensitive to a low fluence and only detects ions.

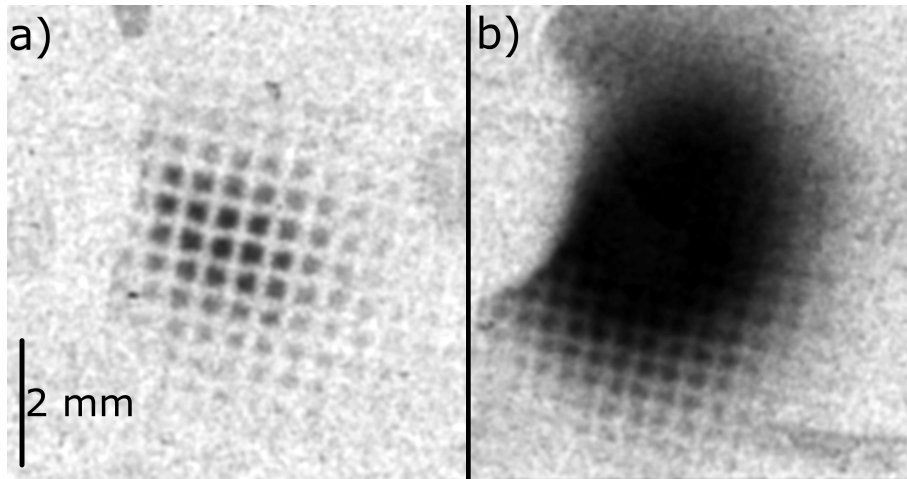


Fig. 5.13: Blast wave proton probing results. a) Shows the mesh imprinted pattern on a scanned piece of etched CR-39 placed 100 mm from the interaction region with average proton energies of ~ 100 keV. The gas jet was not fired for this shot. b) Shows the results with the gas jet firing with CR-39 placed 25 mm from the blast wave. The mesh pattern is distorted and unrecognisable. The round shadow in the left of both images is the gas nozzle with a diameter of 3 mm.

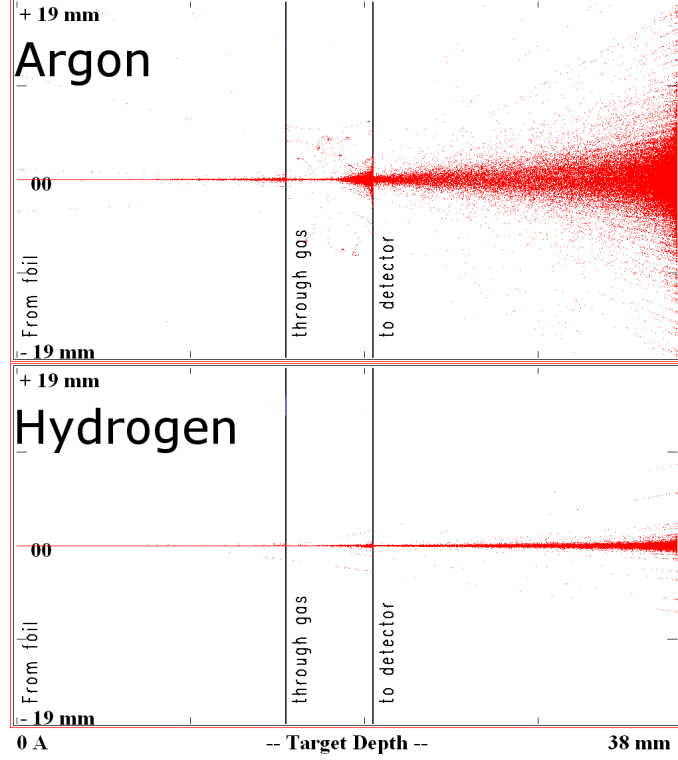


Fig. 5.14: SRIM model output of proton path divergence through argon and hydrogen gas jets. The left and right regions are at vacuum. The gas jet region is at a density of $3 \times 10^{18} \text{ cm}^{-3}$. An initial proton energy of 100 keV is used with no divergence, they finish with an average energy of 50 keV in argon and 82 keV in hydrogen. A clear divergence of $\sim 5 \text{ mm}$ is seen in argon that would easily distort the mesh pattern or indeed any information in the beam.

can penetrate cold cluster gas, but be caught up and compressed by an expanding blast wave and traced via proton imaging will further our understanding of similar dynamics in astrophysical scenarios, for example of the mechanism which allows the solar wind to penetrate the Earth's magnetosphere [137] [138].

5.3 High Energy Scaling of Single Blast Waves

This section reports on observations from the 2007 Vulcan Laser campaign. For a fuller report of the first observation of the thermal cooling instability the reader is recommended towards [28] [6]. This is a significant result as it demonstrates the use of clusters to experimentally demonstrate an effect predicted by theorists to occur in systems such as supernova remnants but not observed until this experiment.

In this section we continue the theme of the chapter and focus on the results that highlight the ability of clusters as a medium to support high energy density physics. This includes their high laser energy absorption even at intensities approaching the relativistic

regime. Also the demonstration of blast wave physics discussed in plasma textbooks but rarely observed except at these high energies such as supercritical shocks, ionisation channels, secondary shocks and surface shock structure.

5.3.1 Experimental Setup

The experiment was similar in concept to the system used in the Glass Laser experimental design (§5.1.2). The green probe beam used Michelson style interferometers instead of Mach-Zehnder. As shown in Fig. 5.15, there were four of these dual, delayed polarisation beams each with a variable delay. One of these four beams was sourced from a Nd:YAG laser instead of the frequency doubled main beam as it was stretched via a multi-element etalon to 30 ns to provide a long pulse for streaking of schlieren images. The main heating beam from the Vulcan laser produced approximately 50 J in 1 ps and was focused using an $f/18$ parabola. A calorimeter was used to record the residual energy of the beam after the interaction. A downward looking Andor camera was used in conjunction with a Celestron imaging lens to image the interaction.

5.3.2 Absorption Results

Previous work by Ditmire *et al.*, taken on an older version of the Glass Laser [127], shown via the blue squares in Fig. 5.16a indicates a drop in absorption above 10^{17} Wcm^{-2} . This was attributed to a *burn through* effect whereby the leading edge of the laser pulse disassembles the clusters before the peak arrives on a picosecond timescale. However, this more recent work carried out at the Vulcan laser extends the laser intensities to 10^{19} Wcm^{-2} and shows that the absorption does not in fact drop. The previously seen drop is tentatively attributed to a large pre-pulse, unique to that laser, disassembling the clusters well ahead of the main pulse.

Fig. 5.16a shows a plateau beyond 10^{15} Wcm^{-2} . This is attributed to a saturation effect whereby no more energy can be effectively coupled into the individual clusters before they explode. Using the downward looking Andor to measure interaction length Fig. 5.16b supports this. Here as individual clusters cannot absorb any more energy there is still significant energy left in the beam to ionise later clusters thereby extending the interaction length as seen. The surprisingly low Z dependence indicates that the saturation depends only weakly on type of gas but more likely on cluster density.

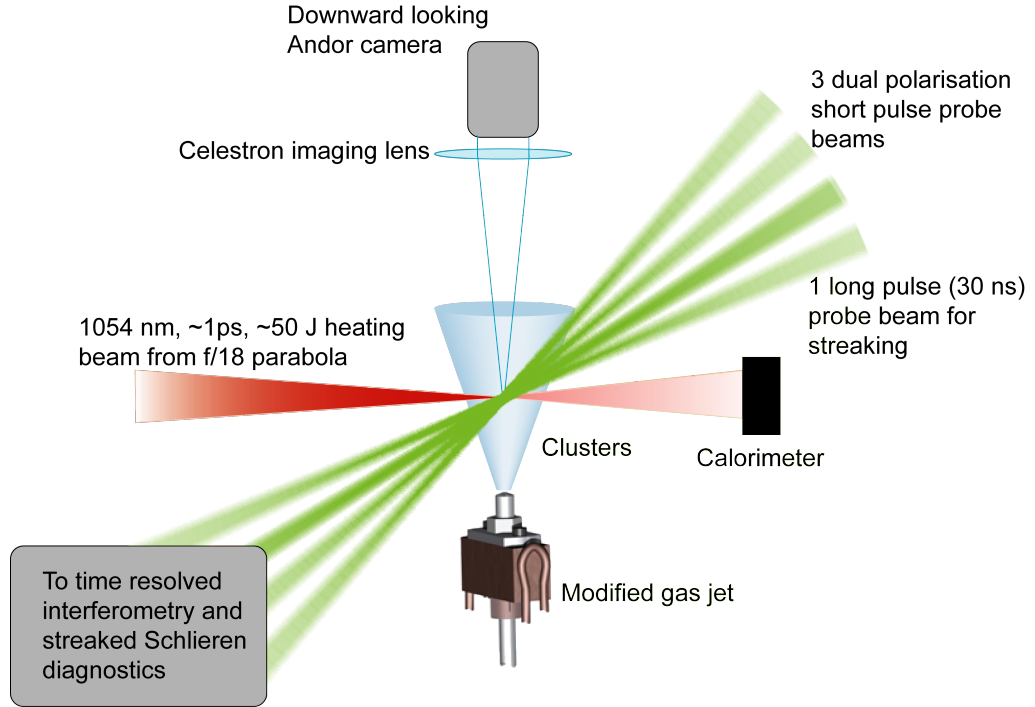


Fig. 5.15: Vulcan 2007 blast wave experimental setup. All four green probe beam shown contained dual polarisations which had a variable time delay as well as their own variable delay as shown in Fig. 5.1. Thus six different timed interferometry and Schlieren images were obtained simultaneously per shot. Geometrically the probes were arranged vertically as close to each other as possible. The long pulse beam for streaking was sourced from a Nd:YAG laser whose pulse had been stretched using a multi-element etalon. Various wavelength filters were used on the downward looking visible Andor in an attempt to separate scatter and plasma self emission. The calorimeter for energy detection was calibrated by the laser staff at the Central Laser Facility. Not shown here are pinhole cameras, single photon Andor detectors for temperature measurements and diodes for measuring side-scattered light.

The low absorption in Fig. 5.16a of hydrogen clusters is not seen elsewhere [127]. We attribute the low (50%) absorption here to the small cluster size (8 nm) and relatively long pulse duration. The nanoplasma model predicts a faster explosion (~ 300 fs) suggesting the cluster won't see all of the ~ 1 ps laser pulse. Also, a reduced resonance is seen for smaller clusters which are much smaller than the skin depth.

In order to calculate the peak energy density a vertically averaged lineout is taken across the downwards looking Andor images, Fig. 5.16c. We see the plateau we would expect from our theory of saturation. A dip occurs directly above the nozzle aperture. This dip may be a result of the laser focus, or even self-focussing, reducing the interaction volume and therefore the deposited energy. Applying the total energy lost, ~ 40 J, to this flat top distribution of length multiplied by the average width, 0.23 mm, for the cylindrical volume we obtain a peak energy density of $\sim 10^5$ Jcm $^{-3}$ in krypton (or ~ 70 GJkg $^{-1}$ which

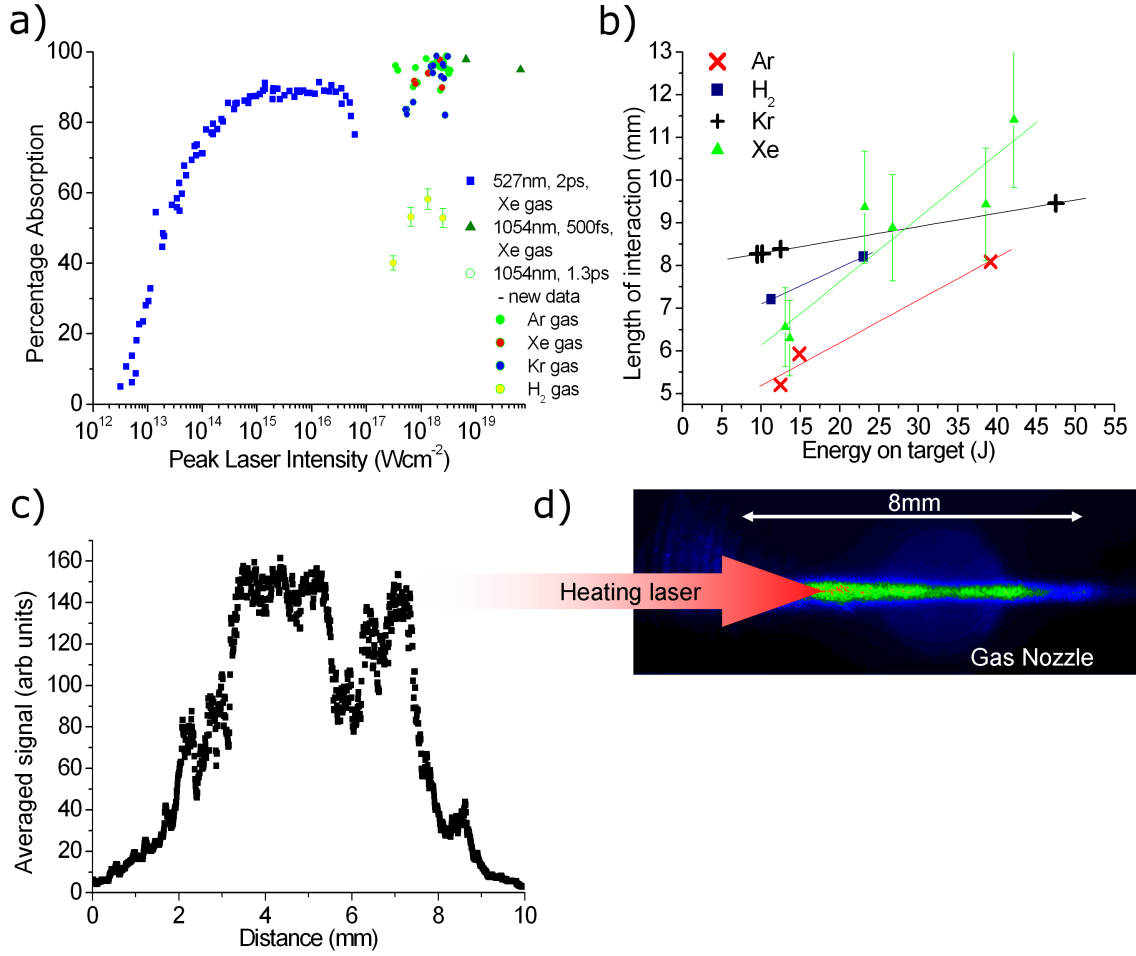


Fig. 5.16: Absorption energy and length data from the 2008 Vulcan TAW experiment. a) Shows percentage energy absorption data. The initial energy is obtained from a calibrated calorimeter behind the last turning mirror provided by laser technicians at the laboratory. The final energy is obtained from a calibrated calorimeter after the interaction. The solid blue squares are reported in Ditmire *et al.* taken on the Imperial College Glass Laser [127]. The solid green triangles were taken on the Vulcan Petawatt laser [139]. The green circles represent new data taken on Vulcan (TAW). b) Shows interaction lengths against energy on target from the downward looking visible Andor. c) Shows a vertically averaged lineout across the Andor image (as seen in d) for a 7.8 J shot in Kr) showing that the energy deposition appears to saturate. The dip in the lineout is directly above the hole in the gas nozzle.

is four orders of magnitude higher than that of TNT). Drake confirms this places us firmly in the high energy density regime [9].

Although our cluster volume as a whole exhibits high laser energy absorption it should be noted that increasing the peak, localised energy density further is non-trivial. Increasing the laser intensity and the cluster density does not increase the energy density. This is a result of refractive defocusing which places a limit on the maximum intensity supported for a given wavelength in a given density. Clusters, which ionise more efficiently than typical

gases and so lead to higher density gradients, may suffer significantly more from this phenomena. Mackinnon *et al.* show at densities of $\sim 10^{19} \text{ cm}^{-3}$ for a $\sim 10^{17} \text{ Wcm}^{-2}$ laser pulse significant refraction begins and energy is lost from the main beam direction [140]. Beyond this the laser light is refracted away from the axial density gradients increasing the effective focal volume. As we achieve higher intensities this becomes more obvious and we see ionisation channels at early time which are also seen by Mackinnon. This is demonstrated in Fig. 5.17 where the channels move away from the high electron density formed in the main heated volume. These plasma channels are thought to arise from refractive splitting of hot spots in the laser beam. Mackinnon tell us that these channels will not in themselves significantly reduce the overall energy deposited in the blast wave however the overall refraction may increase the beams propagating solid angle such that the calorimeter does not measure the whole beam. This may lead to errors in absorption measurements. These channels were mainly seen in argon and hydrogen. They were weakly observed in krypton and not at all in xenon. This is attributed to a reduced density leading to reduced refractive defocusing. This may also account for the generally reduced interaction lengths for argon.

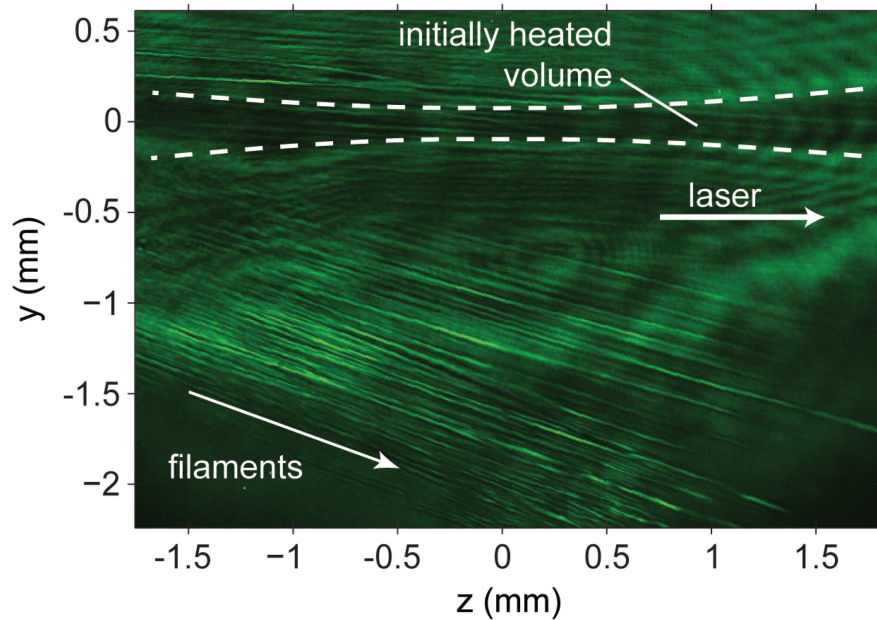


Fig. 5.17: Early time (1.6 ns) Schlieren image of ionisation channels from a 52 J shot in argon from a nozzle backing pressure of 50 bar. The filaments are only seen on the gas jet side of the blast wave where the density is higher. Figure courtesy of Dr Hohenberger [28].

5.3.3 Supercritical Shocks

Zeldovich [21] describes a supercritical shock as one in which the radiative precursor temperature, T_p , rises to approach the value of that behind the shock, T_i . As this occurs the radiation emitted from the T_i region to the T_p region, which is now in near thermodynamic equilibrium, will be re-emitted via radiative heat conduction outwards to a depth depending on the opacity of the ambient material. This leads to a broadening of the T_p precursor region as seen in Fig. 5.18a. The radiative conduction will also mean that more energy is reabsorbed into the shock itself, raising T_s .

In order to observe the supercritical shock the temperature inside the shock, T_i , needs to increase above a critical temperature and as such can only be observed in these high energy experiments. Above this temperature the radiation emitted from the T_i region (approximately black-body radiation as it is in thermal equilibrium) will be high enough to cause a positive net radiation flux for the precursor, raising its temperature to equilibrium with itself. To increase the probability of observing this we can reduce the thickness of the precursor layer that is heated. The opacity of the precursor can be increased thereby decreasing the mean free path of the radiation, or the optical depth. Increasing the gas density will have this effect as opacity is density dependent. This explains why in our data we only see the supercritical shock on the gas jet side of the blast wave where the density is higher by a factor of two.

If the temperature in the precursor increases this should also be observed in the electron density measurements taken on this experiment as increased ionisation will occur.

Fig. 5.18b shows the interferometric analysis of a 25 J shot in krypton clusters held at 42 bar backing pressure and room temperature at 6 ns through the blast wave's evolution. The extended precursor profile which we associate with a supercritical shock can clearly be seen.

Another commonly observed feature in these high energy shocks can be seen to the left of Fig. 5.18b. Here a secondary shock develops from the initial shock front and appears to propagate in the opposite direction through the centre of the blast wave. Previous work at lower energies has shown this feature as a shock front splitting and then slowing [27]. However only at these high energies can it be seen that the secondary shock actively propagates towards and through the central axis [28]. A possible explanation of this is that radiation is emitted from one side of the blast wave and propagates inwards through the low density interior, and therefore low opacity, to impact on the shock front on the

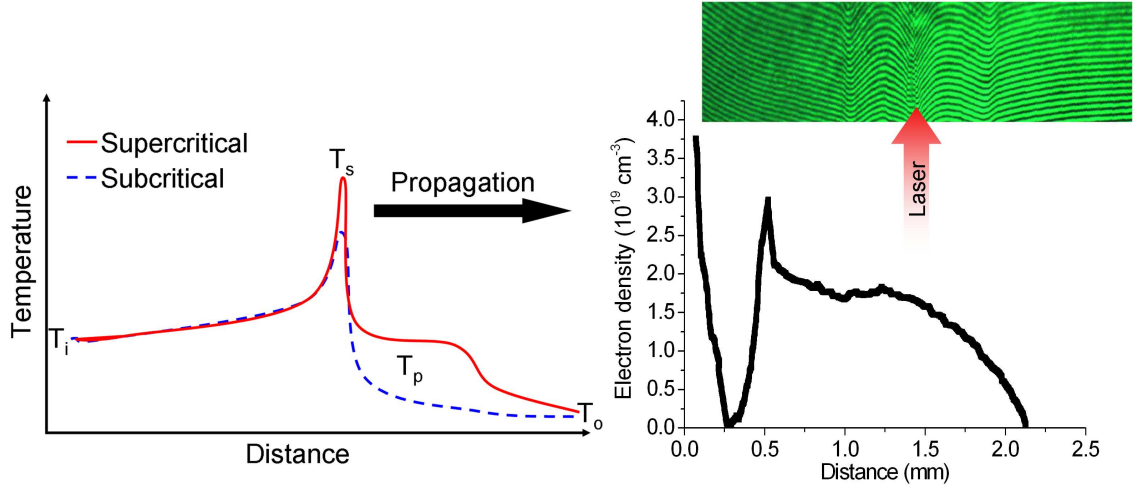


Fig. 5.18: a) Schematic of supercritical and subcritical shock temperatures. T_i represents the temperature inside the shock, T_s the temperature just behind the shock front, T_p the temperature of the radiative precursor and T_o that of the ambient material outside the shock. The supercritical shock has an extended precursor at, or near to, thermal equilibrium with T_i . Adapted from *Zeldovich* [21]. b) Analysis of interferometry (shown inset) from a 25 J shot in krypton. The electron density shape is remarkably similar to the temperature distribution in the schematic. The far left of the graph shows internal shock structure which belongs to another shock travelling internally in the opposite direction. A typical subcritical structure can be seen in Fig. 5.8c.

other side of the cylinder. This could result in an ablation like effect which causes another shock to form propagating inwards and compressing as it travels. Clearly this would only be possible at high energies. Simulations are inconclusive due to a lack of knowledge on opacity and the extent of radiation inside the cylinder. Further work is needed to study this phenomena.

5.3.4 Fine Structure in Shocks

Several of our schlieren images, particularly at high energies, show fine structure which appears to lie on the surface of the shock front. Although these are evident in all four gases we investigated they are most extensive in hydrogen. They also persist to late time, as seen in Fig. 5.19, in hydrogen whereas in the higher Z gases the fine structure smooths out. The fact that they are not resolved on the interferometry images suggests they are on the shock surface not the interior of the blast wave cylinder. All though at first glance this fine structure may seem inconsequential the Schlieren images suggest that these highly localised regions of very high electron density ($>10^{17} \text{ cm}^{-3}$) must have a strong field aligning them in the axial direction and an important underlying physical effect.

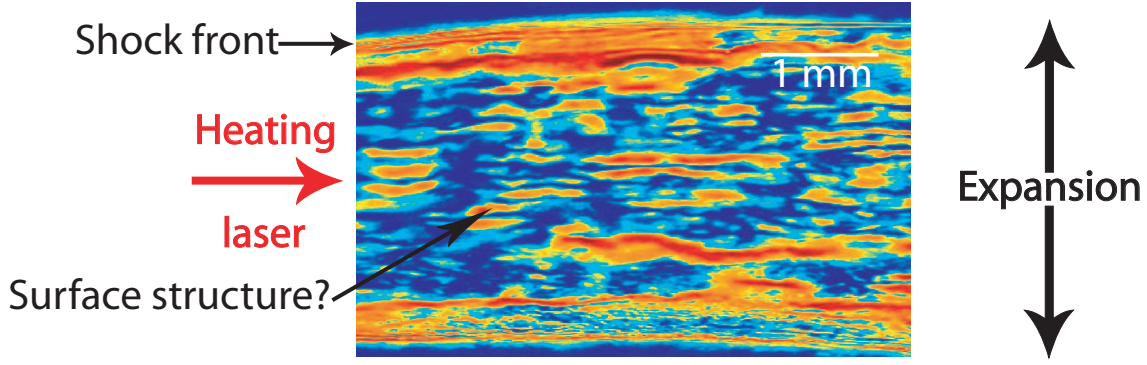


Fig. 5.19: Shock surface structure in a Schlieren image. This image is shown in false colour only to highlight the fine structure. Hydrogen cooled to 70 K at 30 bar backing pressure was used producing clusters with an average radius of ~ 5 nm. Probe delay at 25.4 ns. Elongated millimetre size structures can be seen throughout the blast wave. At the top and bottom of the shock double lines can be observed, this is an artefact of the Schlieren process.

Drake suggests two possible reasons for this fine structure. One option is that they are evidence of an instability causing localised density fluctuations, in particular the *radiative thermal instability* [9]. Here, when material cools by radiating, a temperature instability can develop as a result of the temperature dependence of opacity. A temperature increase will thus lead to less radiative cooling and so to a further temperature increase. This will in turn lead to a density perturbation as the cooler regions are compressed by the hotter ones. We discount this option as hydrogen is essentially non-radiative due to its low Z .

The second possibility is that of an induced magnetic field [9]. Here the space charge separation at the shock front produces a radial electric field which in turn induces an azimuthal magnetic, B , field as given by Ampère's law:

$$\mu_0 I = 2\pi r B \quad (5.1)$$

where Ampère's law is evaluated for an extended wire-like portion of the approximately uniform radial electric field with radius r and current flow I . As a back of the envelope exercise we can see that to produce a significant magnetic field, say 1 GGauss, in a $5 \mu\text{m}$ radius 'wire' with an electron density of 10^{19} cm^{-3} the electrons need to be accelerated to a velocity of $\sim 30 \text{ kms}^{-1}$. This is approximately the same velocity as our high energy blast waves' expansion [28]. Unfortunately, for a uniform static cylinder any such magnetic fields will cancel out since the radial 'wires' exist over the whole surface of the blast wave. However, our cylinder is not static and as well as expanding radially is also expanding axially [121]. This may produce significant transient magnetic fields in the axial direction

where they do not cancel each other out. This would lead to increased localised electron densities as observed. The only method of validating this hypothesis is through laser driven proton probing, providing further motivation for the work presented in section §5.2 of this thesis.

Alex Robinson, a plasma theorist at the Central Laser Facility, carried out magnetohydrodynamic simulations of our blast wave regime with baroclinic magnetic field generation in a 2D cartesian geometry. This magnetic field generation and amplification arises from the electron heat flux [141]. The key aspect of this effect is the transport of a magnetic field to the overdense region by the heat flux and the amplification which is a function of the ratio between convection and diffusion which can result in megagauss fields. The generation is a function of $\nabla n \times \nabla T$.

The simulations show that if there are some strong perturbations to the initial heating pattern, such as may arise from plasma defocusing effects, then magnetic fields do indeed develop, but never to such an extent that the magnetic pressure can rival the thermal pressure. This means that the initial perturbation is gradually smoothed out and the magnetic fields will not alter the density profile. This suggests that while laser driven proton probing may show transient magnetic fields at early time the fields may not be responsible for the fine structure seen here. However, further work is needed as several other magnetic field generating effects exist [141]. Blast waves are well suited for research on these effects as the initial neutral gas allows the generation of external large static magnetic fields which will then be swept up by the blast wave and transported and amplified.

A second feature seen in Fig. 5.19 is a ‘double edge’ to the blast wave at the top and bottom. This is an artefact of the schlieren imaging process and not real. Follow up work has shown that using a horizontal pin as opposed to a vertical pin schlieren stop, while it reduces the sensitivity of the horizontal orientated features, does remove the double edge feature. This effect is attributed to the relatively large rounded edge of the pin stalk causing a diffractive effect at the focus which on the image plane is translated to a horizontal effect. Using a round metallic stop on a microscope slide had partial success at removing this problem however also degraded image quality, as a result of a low definition edge to the stop.

5.4 Summary

In this chapter experimental studies of laser driven high energy density blast waves in a cluster medium were presented. The chapter began with work carried out on the Glass laser at Imperial College using a variable delay, perpendicular blast wave geometry. Interferometric and schlieren probing of these blast waves established that the highly absorbing cluster medium, which can support the initial ionising seed for a blast wave, is destroyed well ahead of the shock and the associated radiative precursor. The nanoplasma model and stopping power analysis of this cluster disassembling, thermal wave shows that it is not the shock front's radiation that destroys the clusters, as originally thought, but ions emitted in the initial laser driven ionisation event. The evidence that the radiative shock is expanding into un-clustered gas supports our use of the thermodynamic properties of standard gas in analysing the blast wave dynamics.

Furthermore, this geometry permitted preliminary experiments in the pre-heating of the upstream gas. This was shown to accelerate the shock front and increase the electron density. This provides the grounding for future more accurate astrophysically relevant work where phenomena such as supernova remnants expanding into lightly ionised matter can be modelled.

The final section presented data observed from scaling single blast waves to higher drive energies on the Vulcan laser at the Central Laser Facility. The ability of the cluster medium to absorb more than 90% of the ~ 100 J laser energy available even at these higher intensities was shown. Although the thermal cooling instability is not presented in this thesis [6], images demonstrating supercritical shocks, ionisation channels, secondary shocks and surface shock structure were shown. These reinforce the aptitude of these cluster experiments for the study of the high energy density physics environment.

The cause of the surface shock structure shown at the end of this chapter was suggested to arise from electromagnetic field fluctuations. This provides another motivation for laser driven proton probing in the future as a reliable means of proving this hypothesis.

Chapter 6

Proton Beam and Stopping Power Characterisation

Chapter 4 describes the development of an accurate method for determining the low energy spectra of a laser driven proton source. The following chapter will start by applying this method for various laser intensities in order to determine how the characteristics of the low energy cut off varies. The work was carried out on the Imperial College Glass Laser and the Central Laser Facility Vulcan Laser. This technique is then developed further to fully characterise the quality of the proton beam by studying the divergence and emittance of the beam with a particular focus on the low energies used in this thesis. This is necessary to evaluate the practicality of using these proton beams as a diagnostic for ionography or properties such as stopping power. Measurements of emittance allow direct comparison with established particle accelerator sources. Also included, for interest, is a brief analysis of the Imperial College MAGPIE pulsed power generator as a source of protons, as this self-emission will give rise to a background signal in future laser driven probing experiments. The Mega Ampere Generator for Plasma Implosion Experiments is used to study wire array Z-pinch implosions and here the author presents the first ion imaging and energy spectra.

Chapters 4 and 5 have emphasised the need for a reliable and accurate method to determine the stopping power of materials to protons. These materials are in the warm dense matter or plasma state and as such are difficult to fully characterise with typical, long bunch length particle accelerator methods. The final portion of this chapter discusses methods for carrying out this measurement out using the low emittance, short pulse proton beams produced in laser interactions.

6.1 Full Proton Beam Characterisation

The following section applies the low energy proton spectra measurement method evolved in chapter 4 at different laser energies. The author believes this is the first time an accurate evaluation of the low energy cut off against laser energy has been carried out. Also, the beam divergence as a function of energy is determined using two different methods. Finally, these measurements are combined to allow calculation of the beam emittance.

6.1.1 Low Energy Proton Spectra at Different Laser Energies

§4.3.2 used a home built magnetic spectrometer (§3.2.3) to measure the low energy spectral profile of a laser driven proton source. Radiochromic film was used in conjunction with this to provide the absolute fluence of this profile. Using this as a benchmark and with knowledge of the beam divergence we can measure absolute fluence using just the spectrometer at a given distance from the source. In other words, knowledge of the solid angle seen by the spectrometer slit for the benchmark fluence allows us to extrapolate the fluence at other distances within that solid angle. It is important to note here that we are looking at the centre of the proton beam. If our sample is larger than the area of the beam the spectrometer samples the spectrum will not be correct.

Fig. 6.1 shows the spectra for four different energy laser pulses and a MAGPIE z-pinch ‘self-emission’ spectrum from a radial wire array geometry shot for comparison. The three higher laser energies were taken on the Vulcan laser (target areas Petawatt and West) with 20 μm gold foils as the proton generating foil. The lower energy, 1 J, was taken on the Imperial Glass laser using a 12.5 μm aluminium foil.

These low energy spectra show two important features. Firstly, all the Vulcan spectra show a low energy peak at 98 ± 2 keV which does not shift with increasing energy. The 1 J data shows a similar peak at 70 keV. This is suggestive about the accelerating mechanism. The limited data indicates a target dependence rather than an energy dependence of the low energy cut off, though more data is needed. The energy independence may be a threshold effect whereby ions with initial energy less than 98 keV can not efficiently exit the rear surface plasma due to electron collisions and as such are not then further accelerated by the leading electrons. The collisions felt by the plasma will arise from the target material’s electrons which will depend on the thickness and type of target material - possibly leading to this target dependence. No current publications can be found in order

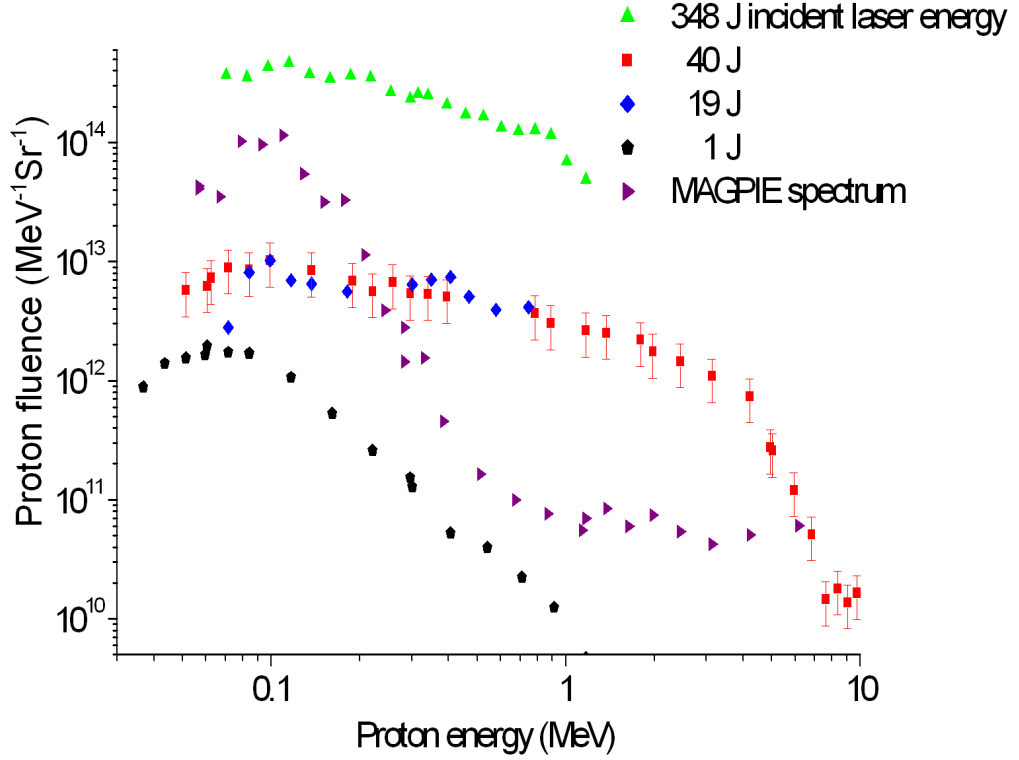


Fig. 6.1: Low energy proton spectra for various laser energies. 348 J, 40 J and 19 J shots were taken on the Vulcan laser (TAW and TAP) with the primary source as a 20 μm gold disc. The 1 J spectrum is an average of 10 shots taken on the Imperial College Glass laser system on to 12.5 μm of aluminium. A MAGPIE proton self-emission spectrum is also shown for comparison.

to provide additional data in this low energy regime for comparison.

Secondly, the overall scaling of fluence with energy is not the established $(I\lambda^2)^{1/2}$ [40] but significantly higher - we do not however state a revised relationship as there are only four data points for comparison. This may indicate that the laser parameters for some of the experiments were not optimal, such as a bad focal spot, large pre-pulse or longer than expected pulse duration. A second possibility is that we are looking at three different regimes: a low energy regime at 1 J where acceleration only just occurs and has not saturated; the 19 J and 40 J regime where the $(I\lambda^2)^{1/2}$ TNSA dependency is typical and the high 348 J energy regime where additional mechanisms may be present such as the RPA process which has an increased energy dependence (§2.2.5).

Clearly more data is needed with more energies, different target materials and thickness. Ideally, this data would be obtained on the same laser, so that laser parameters (other than energy) are consistent. Also, using the same spectrometer to measure the data to ensure consistency at these highly sensitive low energies would be required.

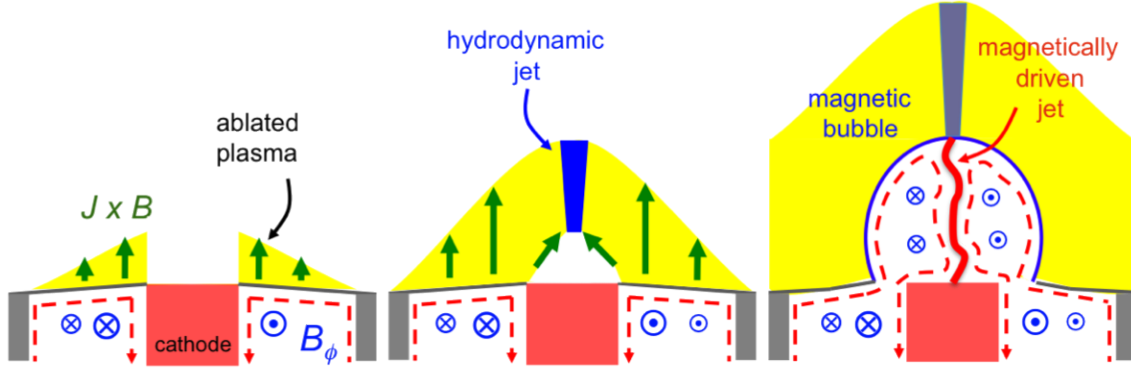


Fig. 6.2: Schematic of a radial wire array experiment on MAGPIE. Current flows radially through 16 10 μm tungsten wires and along the central electrode which drives a toroidal magnetic field below the wires. The plasma ablated from the wires is forced upwards via the $\mathbf{J} \times \mathbf{B}$ force. Complete wire ablation at the central electrode forms a magnetic cavity which is pushed upwards to form a jet by the toroidal magnetic field. Figure courtesy of Dr Francisco Suzuki Vidal.

6.1.2 MAGPIE Proton Self-Emission from a Radial Wire Array

MAGPIE is a 1.4 MA, 150 ns Z-pinch pulse power generator housed in the basement of the Blackett Laboratory, Imperial College. It is used to study the dynamics of wire array implosions and the corresponding plasma features with an astrophysical application [142]. The radial wire array geometry creates a significant magnetic cavity [143] which can launch a magnetically driven jet in a way believed to be analogous to the ‘magnetic tower jet’ produced by a protostar or black hole accretion disk [144] [145]. The physical process to create this magnetic cavity is shown in the schematic Fig. 6.2. The complex magnetic field structure of the jet will be probed by a laser driven proton beam in the near future. As a preliminary to this experiment the proton emission from MAGPIE itself was characterised in experiments with S. Patanker and the MAGPIE team, as this will create a background signal, setting a lower limit for the required energy of the laser driven probe.

The MAGPIE ion energy spectra was measured for the first time by the author using the previously established method with the proton spectrometer at a distance of 16 cm from the interaction centre and can be seen in Fig. 6.1. The maximum fluences captured by the spectrometer are of a level with those observed in the 348 J laser spectra. However, the fluence drops off significantly more rapidly at higher energies such that there are very few protons above 1 MeV. Interestingly, the ~ 100 keV low energy peak is still observed in this data suggesting a broadly similar ionisation and acceleration mechanism to the laser foil interaction may be present.

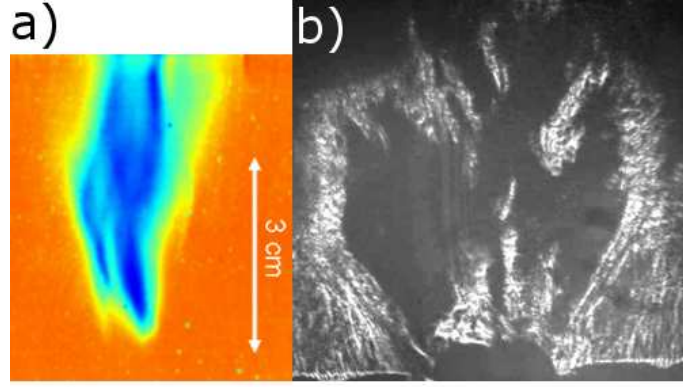


Fig. 6.3: MAGPIE proton pinhole and optical Schlieren images. a) False colour image of a scanned piece of CR-39 from a proton pinhole showing time integrated emission from the MAGPIE interaction. Placed 30 cm from the chamber centre. b) Optical Schlieren image of the interaction at late time ~ 230 ns courtesy of Dr Suzuki-Vidal.

In an effort to spatially localise the emission from MAGPIE a novel proton pinhole camera was developed. This operates on the same principle as its optical equivalent camera using a $100\text{ }\mu\text{m}$ pinhole and CR-39 as the detector. The scanned results along with Schlieren imaging data for comparison can be seen in Fig. 6.3. The emission seems to occur throughout the length of the plasma jet. The emission seen at the top of the plasma jet, which only exists at late time (~ 200 ns), suggests that the protons initially produced may be trapped in the electromagnetic field of the plasma jet, only being emitted as these relax. Fielding several of these diagnostics has shown similar results suggesting a somewhat isotropic emission. This corresponds to a huge overall fluence!

As an order of magnitude estimate of the fluence we take the measured spectra and conservatively extend it over a 2π sr hemisphere. This suggests that the total proton output of MAGPIE is ~ 1.5 J contained in 6×10^{13} protons. As a practical consideration for future proton probing experiments the low energy nature of this background emission will be relatively easy to filter out with aluminium foil when ‘actively’ probing with a 5-10 MeV laser driven source.

6.1.3 Virtual Proton Source Driven by a Low Energy Laser

The TNSA mechanism does not emit protons from the target as a point source. However, for imaging and radiography purposes we can assume quasi-laminar emission and define a virtual source on the laser side of the target, which in turn allows us to define the area of emission from the target [55] [54]. Fig. 6.4 shows the setup used to determine this. A well known mesh is used to imprint a pattern on the proton beam which then expands as per

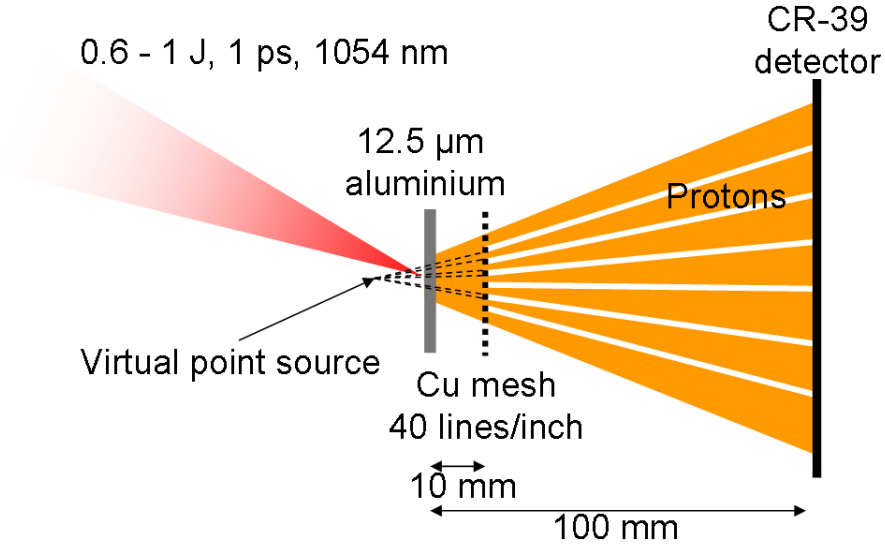


Fig. 6.4: Schematic of virtual source measurement. Using the mesh and the projection of it on the detector backwards ray tracing can be carried out on the individual beamlets to define a virtual point source. Averaging of various energies and mesh positions were carried out to reduce random error. By counting the total number of mesh wires seen the overall divergence from this point source can be established.

the beam divergence on to the CR-39 detector. This can then be ray traced backwards towards the virtual source. We use a copper mesh with a 0.25 mm wire diameter at 40 lines/inch leaving a 0.38 mm nominal aperture - relatively thick wire is used in order to completely stop the protons, only energies above $\sim 10 \text{ MeV}$ will penetrate this thickness. Averages of ray tracing of both the wire shadows and aperture beamlets were taken to reduce any systematic error from particle scattering. In order to reduce the random error, particularly introduced in the measurements of mesh, target and detector separation ($\pm 0.5 \text{ mm}$), the experiment was repeated at different energies ($0.6 - 1 \text{ J}$ for the Imperial College experiments) and mesh distances and an average taken. Typical distances are shown in Fig. 6.4.

An average virtual source distance of $650 \pm 200 \text{ }\mu\text{m}$ in front of the target is seen. This is a factor of three higher than seen in other work at higher intensities [55] [54]. This suggests either a high energy dependence or, more likely, that the experiment is influenced by an electric charging of the mesh which alters the divergence mid-propagation. Indeed, not grounding our mesh resulted in substantially fewer protons detected. At our substantially lower laser energies this effect would be dramatically reduced.

By counting the number of wires imaged on the detector we can establish the maximum angle of divergence from this virtual source and hence the area of emission on the real

target. Our half angle of divergence is $18 \pm 3^\circ$ where the large error arises from the relatively large mesh size. This gives us an area of emission with a diameter of $422 \pm 150 \mu\text{m}$. Again, this is a factor of two larger than other work. We suggest this is an effect of the lower laser energy operating in a different, unsaturated regime. Here, the lower overall energy of the protons produces a slower breakout from the plasma allowing more time for energy transfer and further ionisation of surface contaminants. Alternatively, this may be a symptom of a larger than expected focal spot. Further work is needed at these low energies which may push forward understanding of the TNSA process.

6.1.4 Energy Dependent Divergence

The nature of the high energy proton beam on the warm dense matter Vulcan laser experiment permitted the use of a radiochromic film detector stack. Here, the energy selectivity of each layer allows measurement of the energy dependence of the angular divergence. Fig. 3.14 shows us that for each layer of the stack the dose is dominated by the protons which Bragg peak in that layer and as such we can ignore the dose deposited by higher energy protons with a relatively small increase in error. The trend in Fig. 6.5 shows that divergence decreases as energy increases. This means that neglecting the removal of dose contribution in a layer from protons which Bragg peak in later layers will not effect the dose on the outside of the beam - the region used for divergence analysis.

Fig. 6.5a shows lineouts taken from an RCF stack used in the warm dense matter experiment on the Vulcan laser. The RCF was scanned using a standard flatbed scanner and the pixel values converted to dose as detailed in chapter 3. The line feature discussed in §4.3.2 can clearly be seen at -11° and may alter the divergence. Generally the profile can be seen to be approximately flat-top rather than Gaussian. Layers 1 and 2 are beyond the dynamic range of the scanner. Layer 1 may have contributions to the dose from other radiation, for instance heavy ions which penetrate the $12.5 \mu\text{m}$ of aluminium on the front of the stack. Fig. 6.5c shows the maximum half angle of divergence till the dose drops to 10 Gy, chosen to be clearly above the noise threshold. This data is averaged over three stacks and several lineouts. The large error bars arise from both a lack of symmetry in the beam and shot to shot fluctuations.

The half angle of divergence is seen to vary approximately linearly with energy which agrees with other work reported in the literature [50]. However, the low energy divergence is significantly higher than expected and does not feature the plateau seen in other work.

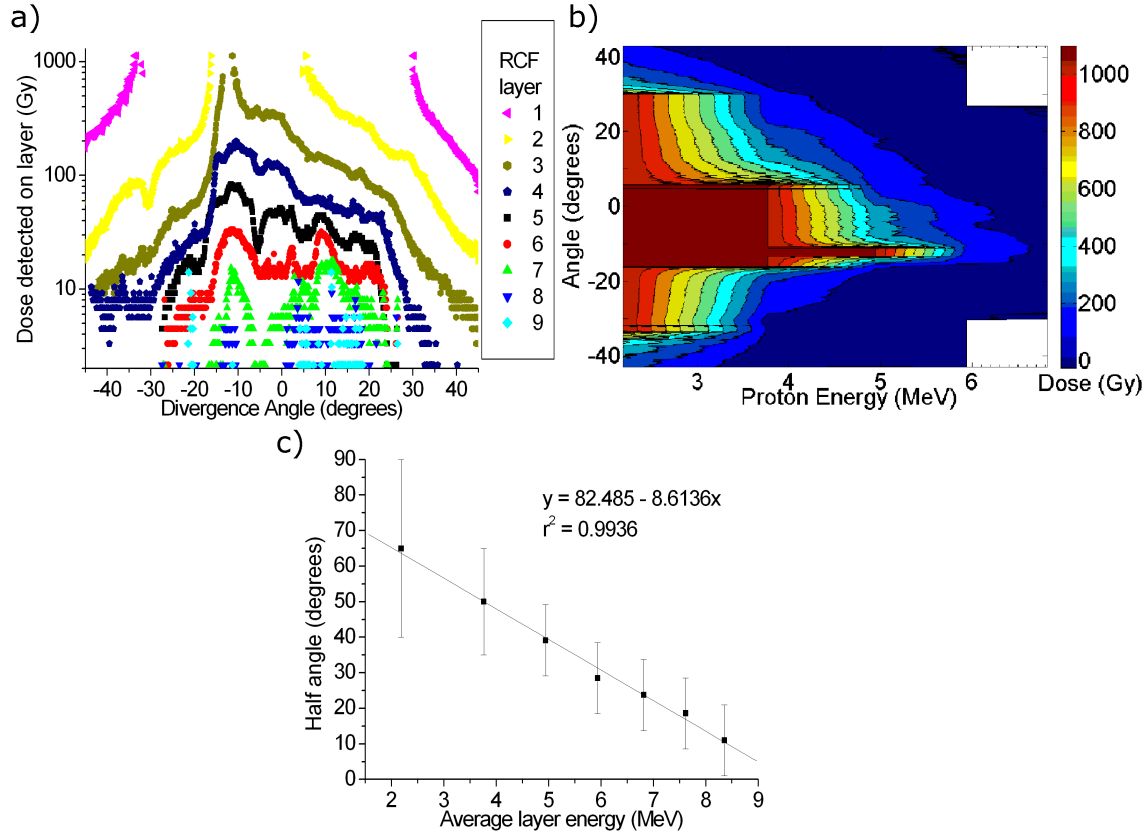


Fig. 6.5: Proton energy dependence of angular divergence from a 79 J (pre-compressor) shot into a 20 μm thick gold disk on the Vulcan laser. a) Lineouts of standard flat bed scanner images with pixel values translated to dose for layers 1-9. Lineouts of layers 1 and 2 are saturated. The peak at -11° is a result of the line in the RCF, seen in Fig. 4.9b, and is omitted for maximum divergence measurements. b) Contour plot showing the same data as a). The saturated region can clearly be seen in dark red at low energies. c) Shows half the divergence angle envelope for each layer's dose to drop to a 10 Gy threshold. Errors arise from shot to shot energy fluctuations and the lack of symmetry in the beam profile. Layers 1 and 2 are extrapolated and as such have a larger error. Layer 1 may also have contributions from other radiation which was not adequately filtered so may be an overestimate.

While much of this may be attributed to the large error a possible physical cause may be the small, 200 μm , diameter of the proton producing foil. Here plasma formed on the rear of the foil may 'leak' around to that formed on the front and alter the acceleration dynamics for the lower energy protons which exist on the plasma edges. Other explanations exist such as a curved target which acts to defocus the beam or a large laser pre-pulse which may form a curved target. These differences show that it is important to characterise the proton beam for every experiment where a different laser and target is used. Within these experiments shot to shot variations are minor.

6.1.5 Low Energy Transverse Emittance

In the previous two sections a laminar proton beam has been assumed. A laminar beam is one where all particles at a position have identical transverse velocities and the magnitude of this velocity is linearly proportional to the displacement from the central axis [146]. In other words, the particles that start at a point source will never cross trajectories. If a beam is not laminar then the point projection model used in the last two sections will be inaccurate as, for example, the beam may suffer from charge repulsion. To measure the degree of laminarity the variable *emittance*, ϵ , is introduced. The smaller the emittance the better the ‘quality’ of the beam. One may expect that low energy, ~ 100 keV, beams such as those produced by the 1 J Glass laser have an increased emittance as they travel slower and are more susceptible to transverse forces driven by space charge.

Ion and electron beams can be fully characterised using a six-dimensional phase space, typically consisting of the three Cartesian coordinates and their associated velocity or momentum (x , y , z , ν_x , ν_y and ν_z). The emittance is the area of a bounding ellipse on a plane connecting data from any two of these six properties. Here the particle accelerator community convention is followed and we define transverse axes x and x' where x' is the transverse angle a particle makes with the central z -axis ($x' = \tan(\nu_x/\nu_z)$) [146]. Liouville’s theorem states that the emittance is conserved when no external forces or apertures act on the beam, hence measurement at the detector should equal that at any point along the propagation in free space. Acceleration reduces emittance as the transverse angle will be reduced over a given distance. As such the *normalised emittance*, ϵ_n , that remains constant during acceleration is used. For a non-relativistic beam the normalised transverse emittance over x is $\epsilon_{nx} = \epsilon_x \beta = \epsilon_x \nu / c$ [146]. The relativistic beam multiplies by a γ factor. For a symmetric beam ϵ_x should be equal to ϵ_y . The longitudinal emittance, ϵ_z , is more important in higher energy beams where energy divergence occurs.

Fig. 6.6 shows the data obtained from a 0.6 J shot on the Glass laser. Using the spectrum in Fig. 6.1 the average proton energy is ~ 100 keV which gives a β of 0.015. Averaging over several shots and axes gives a transverse normalised emittance of 0.07 ± 0.02 $\pi \cdot \text{mm} \cdot \text{mrad}$ at 100 keV. This compares well with the CERN Super Proton Synchrotron which has a rms normalised emittance of $0.4 \pi \cdot \text{mm} \cdot \text{mrad}$ with 10^{11} protons per bunch over nanosecond time-scales. Although, the synchrotron is capable of producing monoenergetic 450 GeV protons bunches! The lower energy seems to have little effect on the emittance in comparison with other work at higher energies [54] [55].

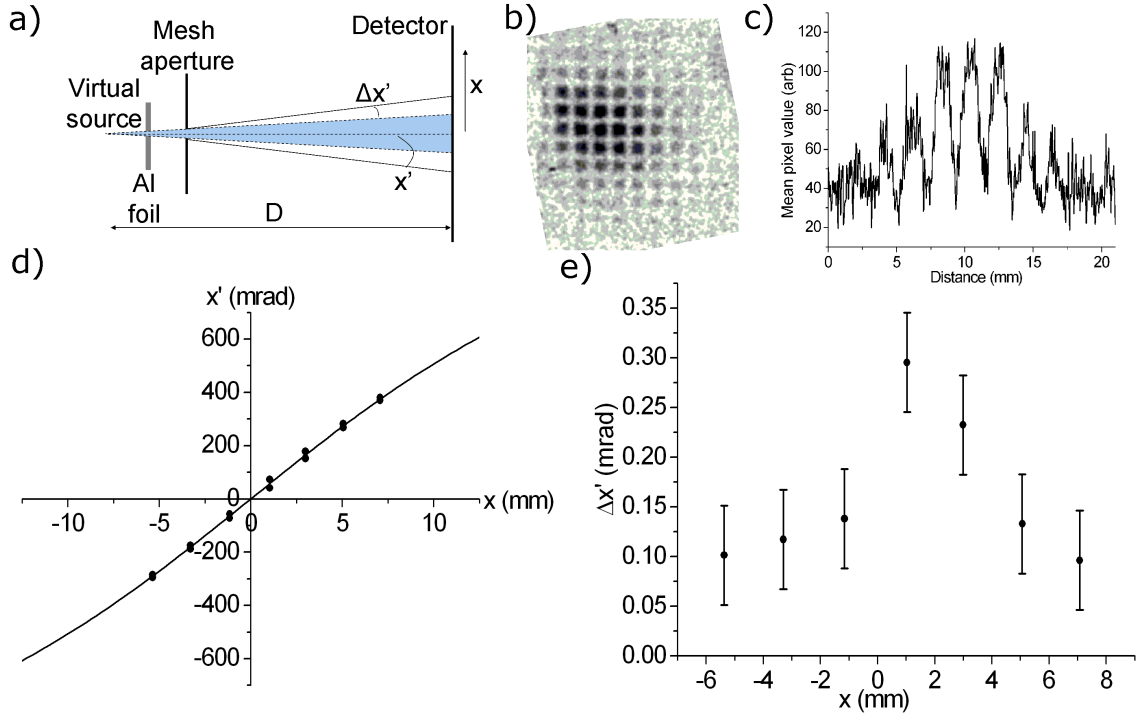


Fig. 6.6: Schematic and results of low energy emittance measurement. a) Schematic of beam as it travels through a mesh aperture. The nominal geometric solid angle is shown shaded. If the beam blooms or the virtual source has a finite size the solid angle will enlarge by an angle $\Delta x'$. b) Etched CR-39 detector showing protons signal with the mesh pattern imprinted for a 0.6 J shot with the mesh 5 mm and the detector 18 mm from the source. c) Lineout of a column of holes averaged across their width to reduce noise. The width was determined at 90% of the peak value. d) Transverse phase space plot of the proton beam. The solid line indicates a theoretical purely laminar beam with purely geometric divergence. The error bars are given by the size of the data points. The area of an ellipse containing all the data points over the extent of the black line gives the emittance. e) Reduced x' showing the blooming away from the theoretical purely laminar case. The large error bars arise from difficulty determining the width of the signal due to noise and show the necessity of averaging over a large data set.

For practical projection purposes we can describe the non-laminarity of the beam as the virtual source having a finite size such that the protons are emitted from an area. The maximum size of this area can be established via back tracing the line marking $\Delta x'$ in Fig. 6.6a for the maximum value of $\Delta x'$. This gives a maximum diameter of $3.2 \pm 0.4 \mu\text{m}$.

Knowledge of the emittance now allows us to deduce the error in measurements of density, electromagnetic fields or stopping power where a ballistic geometric trajectory is assumed. However, this low value of emittance means this assumption of laminarity will induce very little error.

6.2 Determining the Stopping Power of Warm Dense Matter and Plasma

Chapters 4 and 5 showed the necessity for a good knowledge of a material's stopping power for ions either to evaluate energy deposited in a material or the velocity of the ions as a function of distance. For a broad range of condensed cold matter the stopping powers are well known and can be found in codes such as SRIM. However, for the warm dense matter and plasma states the stopping power is not known. A literature search shows very little experimental work has been carried out [147]. Several theoretical models have emerged recently, motivated by the possibility of using laser generated ions for the heating stage in inertial confinement fusion [148].

The dearth of experimental data arises from the transitory nature of these states in the laboratory and the difficulty in creating them. Typical condensed matter stopping power characterisation methods involving particle accelerators [149] are not practical for current WDM sources as the hot dense matter does not exist in a steady stable state in the laboratory beyond the nanosecond time-scale - an ultra-fast source of probe ions is thus required. Also, before the stopping power can be determined the hot state itself needs to be created and characterised, the methods for which are still being evolved. Ultimately, once the state has been created, simultaneous time resolved characterisation of the material properties and the stopping power need to be carried out such that the stopping power can be assigned to specific temperatures and densities as the material evolves. Measurement of the initial and final ion energy, E_i and E_f , will give the stopping power through a material of distance z :

$$-\left. \frac{dE}{dz} \right|_{E=E_i} = -\lim_{\Delta z \rightarrow 0} (E_f - E_i) / \Delta z. \quad (6.1)$$

Experimentally the averaged stopping power is measured:

$$-\left.\frac{dE}{dz}\right|_{E=E_f-\frac{\Delta E}{2}} \approx \frac{E_i - E_f}{\Delta z}. \quad (6.2)$$

The stopping power varies at different ion energies as slower ions are more vulnerable to nuclear stopping. The electronic stopping power will peak when the ion velocity approximates the mean velocity of the target electrons. Inside the plasma Debye sphere the ion undergoes typical free electron stopping, the energy is then transferred to electrons outside the sphere via collective excitations, plasmons. To ensure the stopping power is accurate and confined to one of these regimes ΔE should be small - typically $\Delta E < 0.1E_i$. This sets limits on the sample size or ion energy.

Chapter 4 described a method for creating and characterising hot dense matter via proton heating and a method for producing the short, 1 ps pulse of ions required to probe the state. Hence, with two $\geq 10^{20}$ Wcm⁻² short pulse lasers, such as on the AWE Orion laser, fuller characterisation of the material's stopping power, $S(E, T(t), \rho(t))$, should be possible in the future. Ideally a tunable mono-energetic ion beam would be used to enable direct E_i to E_f mapping over a range of energies. However, no current method exists to directly generate these with lasers. A method using the current generated spectrum, such as in Fig. 4.8, is necessary.

6.2.1 Methods of Characterisation Using Laser Generated Protons

To determine stopping power it is necessary to accurately measure the mapping of a specific input energy to a specific output energy, preferably within the condition that $\Delta E < 0.1E_i$. This provides a definite limit for the minimum useful input energy, E_{min} , such that it is significantly above the cut-off energy, E_{cut} , where the protons do not penetrate the material and are fully stopped.

To develop the method it is useful to work initially with condensed matter, such as aluminium, where the stopping powers are well known and are shown in Fig. 6.7. Fig. 6.8 shows the proton spectrum measured in Chapter 4 with the theoretical spectrum after 6 μ m of aluminium calculated using SRIM with E_{min} and E_{cut} marked. It can clearly be seen that the spectrum limits the maximum size of sample - in order to get a large useful region we need a thinner sample. For example 1 μ m of aluminium gives an E_{min} of ~ 600 keV. Ideally, the useful region will fall in the area of greatest dispersion, or maximum

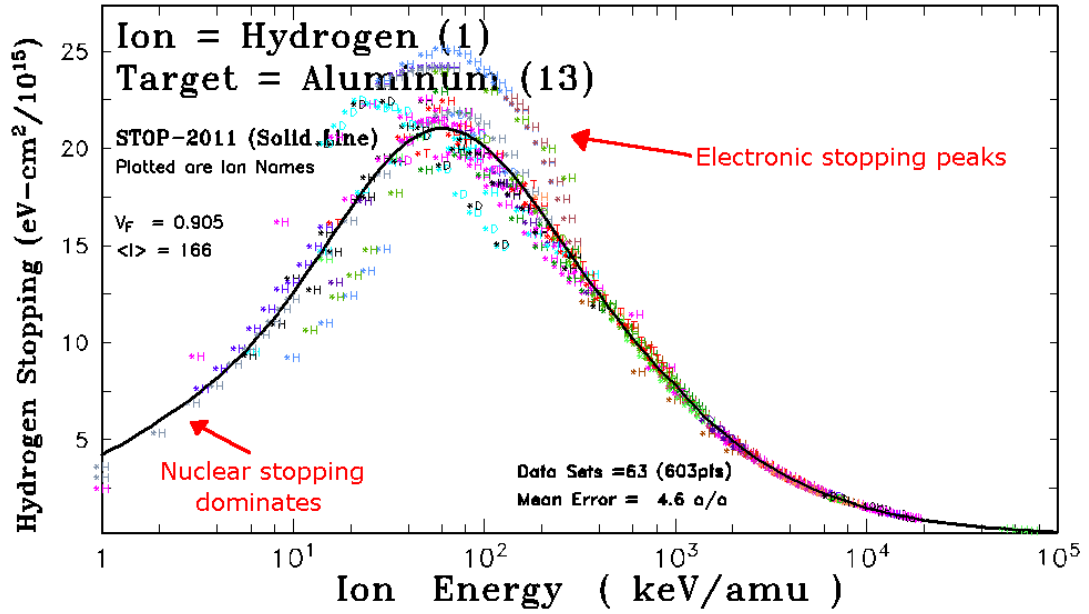


Fig. 6.7: Experimental data SRIM uses for cold aluminium stopping power. Adapted from the SRIM user manual. The data points come from 63 separate experiments (see SRIM for citations). The solid black line shows the semi-empirical model SRIM uses. Such a multitude of data does not currently exist for the hot dense state of aluminium. Though we expect a similar curve for our measurements.

sensitivity, of the magnetic spectrometer.

Having established the useful region for stopping power determination Fig. 6.8 suggests that to map a specific E_i to a E_f we simply pick a point on the input spectrum and move horizontally along the axes till we meet the output spectrum at the same fluence. However, the first attempt at this method highlighted an obvious problem - the fluence does not stay constant on the detector.

The experiment was carried out using the previously described magnetic spectrometer with 6 μm aluminium foil covering half of the entrance slit. This method was used in order to sample the two spectra as close together as possible - in other words the foil stops the same incident beam that the unfiltered half of the spectrometer detects. Using 250 μm slits and a 40 J Vulcan laser shot as expected significant signal was detected on the unfiltered side. However, a factor of four less fluence was seen on the filtered half - barely above the noise threshold. As seen in Fig. 5.14 a proton source will undergo small angle scattering such that it appears to diverge as it propagates through the sample - less fluence will reach the detector plane of the magnetic spectrometer. This also has the side effect of decreasing the energy resolution as protons are scattered to appear at different places on the energy detection plane. To reduce this effect the sample should be placed far away

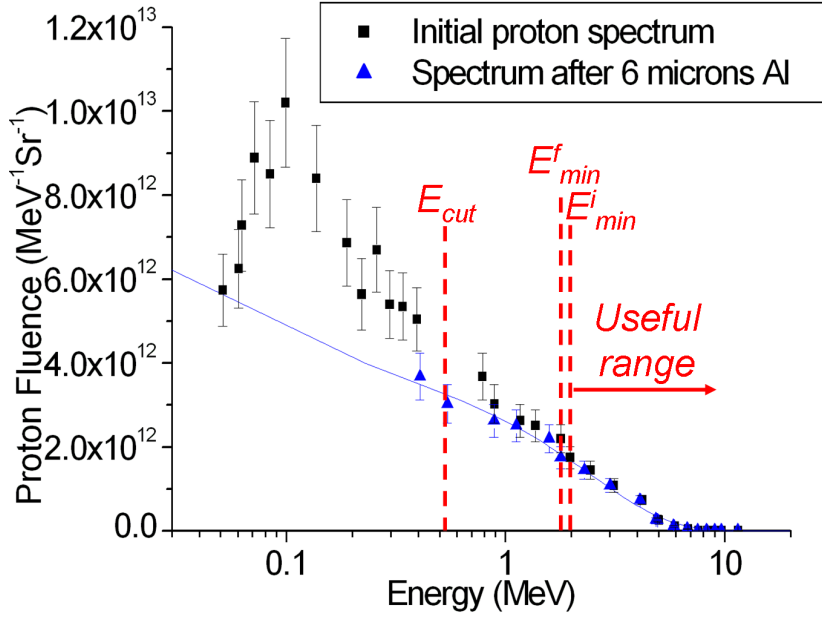


Fig. 6.8: Schematic showing the useful range for stopping power determination of 6 μm aluminium. The black square data points show experimental spectra measured in Chapter 4 from the Vulcan laser. Blue line and triangular points show the same spectrum theoretically after 6 μm of aluminium modelled using SRIM. E_{cut} shows the threshold input energy below which protons are completely stopped in the sample. E_{min} shows the energy above which $\Delta E < 0.1E_i$.

from the double entrance slits however this further reduces the fluence. Unfortunately, the scattered protons will also fall on the unfiltered half of the spectrum and erroneously increase the incident spectra fluence.

To overcome these difficulties a separating plate can be used between the two spectra - from the sample all the way to the detector plane. Also, a divided slit with the sample side considerably wider than the unfiltered side. Although it may be easier to use two spectrometers this method means that the E_i and E_f spectra are measured as close together as possible and with the same calibrated magnetic spectrometer thereby reducing error. Unfortunately, laser time constraints did not allow the application of these modification. In an attempt to approximate this method we can compare a ‘standard’ measured proton spectrum with one taken on the next laser shot under the same conditions with foil over the whole double slits which were held at 1 mm wide.

Using this *quasi-two spectrometer* method it is possible to obtain an E_i and an E_f spectrum on one shot. However, we cannot directly compare their fluences to map a specific input energy to a specific output energy due to scattering effects. It would be necessary to use a large area RCF to capture all the scattered protons and attempt to

recalibrate the spectrometer for the filtered spectrum and adjust for the scattering effect. This would have to be carried out on every sample and, particularly for the time varying hot dense states, is not practical.

An alternative method uses the gradient of the spectrum as opposed to the absolute fluence. Here, a sharp feature of the spectrum shape in the input beam is mapped to a similar feature in the transmitted beam. For our spectra this would be the low energy peak at ~ 100 keV. To place the low energy peak in the useful range it would be necessary to use a 80 nm thick piece of aluminium, though this width may be significantly thicker for low density plasma targets. This has the advantage of being in the most sensitive region of the spectrometer. Unfortunately, as established in §6.1.1 the low energy peak cannot be shifted to higher energies so this method can only be used at $E_i \approx 100$ keV. Other work shows different structures in the spectra at higher energies, which potentially allows thicker targets to be studied, via the use of target heating, pre-pulse or the use of the radiation pressure acceleration mechanism.

An approximate solution to the stopping power determination at higher energies is to use the low energy peak and other significant features method to calibrate the transmitted spectral fluence. It was established earlier that if not for the scattering effect we could directly map E_i to E_f by looking at points on the two spectra with the same fluence. By this argument the low energy peak energies on the two spectra would be at a constant fluence, merely shifted in energy, if not for scattering effects. Hence, we can lift the transmitted spectra up so this is true. Unfortunately, this is not sufficiently accurate as the scattering effects are energy dependant and it would be necessary to use several of these significant spectral features and extrapolate an energy dependent scattering fluence correction. This will of course have significant error but if it produces a fluence calibrated transmitted curve will allow stopping power measurements over a large range of energies. Also, total deposited energy calculations where this is equal to the difference in the area underneath the two spectra.

Clearly using the laser generated proton spectrum to capture stopping powers over a range of energies in one shot is non-trivial.

6.2.2 Methods Utilising Mono-Energetic Proton Beams

The ideal solution for stopping power analysis is to use a mono-energetic beam as this allows direct E_i to E_f mapping. It is also necessary to have an ultra-fast proton beam

to resolve the stopping power at specific transitory temperature and density states. The author suggests a method to combine these two requirements. Fig. 6.9 shows the experimental concept. Here, two laser driven proton beams are generated for use in the pump probe experiment. The heating proton beam passes through a heavy ion filter and goes on to heat the sample whose temperature and density state is diagnosed as described in chapter 4. The proton probe beam is generated from a curved target in an effort to reduce the angle of divergence and increase fluence on the WDM sample. If necessary an Einzel lens can be used to focus the beam further, selecting focussing voltages suitable to the desired proton energy [150]. A set of curved conducting plates with opposing voltages is then used to select the desired proton energy. This electromagnetic selector is commonly known as an electrostatic analyser. The protons energies not selected will not follow the physical geometry of the plates and collide with them. This picosecond quasi mono-energetic proton probe will then traverse the sample and an energy change will be detected on the magnetic spectrometer. Commercial code, such as SIMION [151], exists to carry out ion tracing so that the correct voltages can be selected.

This method will be equally applicable to diagnose the stopping power of plasma where the sample wire is replaced with a plasma created through direct laser heating. The disadvantage of this method compared to using the whole proton spectrum is that only a single energy stopping power value is measured per shot. Thus many shots need to be carried out to fully diagnose the stopping power for even a single temperature and density state. With the high power laser systems' low repetition rates this will take an extended period of time. Generating the necessary large voltage for the energy selection may also be problematic.

One of the requirements for this method, particularly in probing warm dense matter, is a very high intensity laser $> 10^{20} \text{ Wcm}^{-2}$. The strength of the warm dense matter heating method described in chapter 4 is that it heats microscopically large samples of 5-10 μm in size. Thus, the useful range of protons is at high energies as seen in Fig. 6.8. Unfortunately, the fluence drops off rapidly at these high energies. Hence, to produce the required signal and lengthen the energy range probed a very high intensity laser is necessary to generate a proton beam with an increased high energy cut-off and fluence.

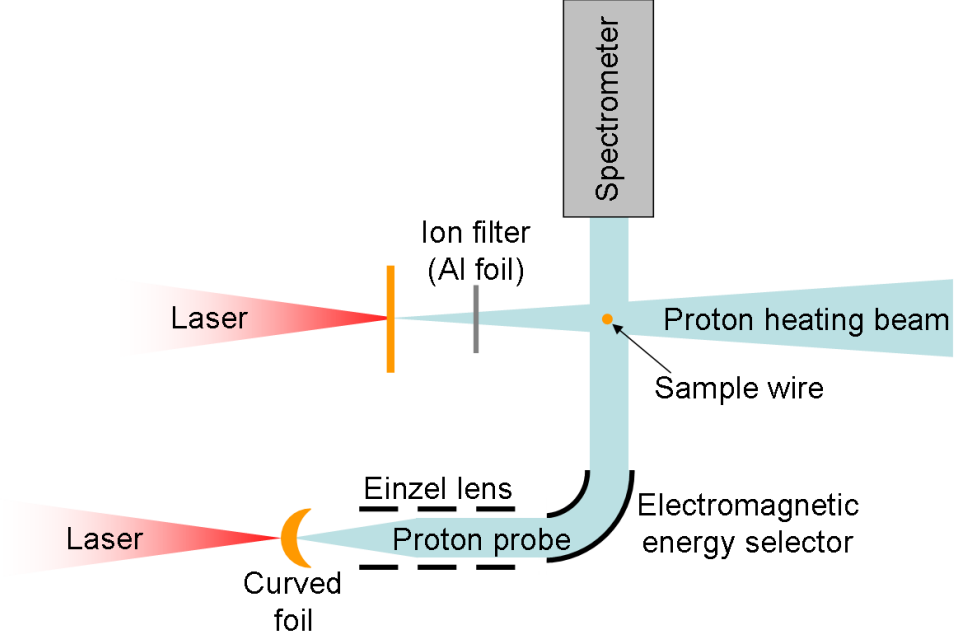


Fig. 6.9: Schematic for a quasi mono-energetic laser driven proton probing beam. Not shown here are the diagnostics to characterise the sample temperature and density state. A curved generating foil is used to produce a proton beam with a reduced half-angle of divergence, increasing the signal reaching the sample. The desired energy is selected by using a curved set of plates, commonly known as an electrostatic analyser. Here, the protons are bent via the Lorentz force as a specific voltage is applied to the plates. Only a narrow range of energies will exit the plates without hitting them. Hence a near mono-energetic beam is created of picosecond duration. This probe beam will then pass through the heated sample and be slowed by the stopping power of the sample. This new transmitted energy can then be measured on the magnetic spectrometer. If necessary to increase the signal an Einzel lens can be used to focus the protons with the desired energy onto the the entrance aperture of the curved plates.

6.3 Summary

In this chapter further experimental characterisation of the laser generated proton beams used in chapters 4 and 5 were presented. This included both the low energy proton beams generated from the Imperial College Glass Laser and the higher energy beams generated from the Vulcan Laser (Target Area West) at the Central Laser Facility. For the first time the accurate low energy end of the proton spectrum (70-1000 KeV) is presented at varying incident laser energies (1-348 J). The scaling does not follow the established $(I\lambda^2)^{1/2}$. As an aside, work carried out diagnosing the proton emission from the Imperial Z-pinch, MAGPIE, was shown as a prelude to laser driven proton probing of a magnetic tower jet experiment. Then for projection purposes the virtual point source was introduced as a method of determining the area on the foil which emits protons - $422 \pm 150 \mu\text{m}$ diameter on the ~ 1 J Glass laser. The radiochromic film stacks from the warm dense matter Vulcan

experiment showed that the proton angle of divergence is linearly dependant on energy. The final proton beam characterisation section measured the transverse emittance of the low energy proton beam produced by the Glass laser. This was found to be similar to that observed at much higher laser energies. The very low value of $0.07 \pm 0.02 \pi \cdot \text{mm} \cdot \text{mrad}$ at 100 keV implies a very high quality beam, with near laminarity which is required for the use of the beam as an imaging diagnostic, in for example cluster blast wave studies.

The second half of the chapter focussed on methods to determine the stopping power of warm dense matter and plasma states. The problems associated with using the whole laser generated proton spectrum to determine the stopping power at a range of energies in one shot were described. Although this method is ‘possible’ it is in many ways impractical to implement well. Instead, a method which uses the advantage of the short pulse length of laser generated protons is put forward where an electrostatic analyser is used to select a narrow energy range from the spectrum where the energy lost due to the stopping power of the sample can be directly mapped. A disadvantage of this method is that many shots are required to build up an energy dependant stopping power curve for even one temperature and density state.

Chapter 7

Conclusions and Future Work

In this thesis, we have discussed experimental laser-driven studies of warm dense matter and higher energy density plasmas with a specific focus on laser driven ions as a tool for both heating and measurement. Chapter 1 provided introductory material and described the thesis structure and contributions from the author. Chapter 2 began with a review of plasma physics and an introduction to the warm dense matter regime. The comparatively long range of laser-foil interaction generated ions in material and their short bunch length currently makes them the best candidate for heating a microscopically large sample rapidly to a single temperature and density state in the warm dense matter regime. The highly efficient absorption of laser energy by atomic clusters also enables us to go further via direct laser heating and reach the high energy density plasma regime. The chapter concludes with basic shock theory for blast waves formed in these cluster plasmas and their suitability for scaling to provide a comparison with astrophysical systems such as supernova remnants.

The third chapter defined the instrumentation used in the later experimental chapters. Starting with the non-linear optics theory underpinning the design of the Imperial College Glass Laser and the Rutherford Appleton Laboratories Vulcan Laser systems. Optical probing techniques including interferometry, Schlieren and streaking were described as a method for recovering electron density values in both space and time. Particle probing methods utilising ions to detect both density and electromagnetic fields were introduced. This led to a report on the detailed calibration of a magnetic spectrometer and radiochromic film for low energy protons using a particle accelerator source. Finally, the production of atomic clusters via a modified, cooled, pulsed solenoid gas-jet was described.

Chapter 4, the first fully experimental chapter, contained warm dense matter research begun at the AWE Helen laser in 2008 and continued on the Vulcan laser (Target Area

West) in 2010. Several types of wire samples were heated via a laser produced proton beam. The density was recovered from the Abel inverted transmission of an XUV back-lighter which was streaked over time. The Foord theoretical method allowed pressure recovery under certain conditions and assumptions. Near uniform heating was observed. The UV self-emission from the wire was also streaked with only partial success due to a low signal to noise ratio. The author fielded diagnostics to resolve accurately low energy proton spectrum. This enabled calculation of the deposited energy using the cold stopping power code SRIM. This calculation agreed broadly with the theoretical deposited energy based on the density expansion measurements. The differences between code and experiment highlight the need to properly measure stopping powers in the warm dense matter regime. While the experiment succeeded in creating and characterising the warm dense matter state, improvements will be made on future AWE Orion experiments to ensure more uniform heating and fuller characterisation of the ionisation and temperature states accessed over time.

Chapter 5 experimentally develops previous work on laser driven blast waves, particularly with reference to their applicability for comparison with their astrophysical analogues. As the blast wave expands the interaction with upstream matter affects the dynamics. Clustered gas may offer, for example, a higher opacity to a radiative blast wave than a typical monatomic or molecular gas. Using a new two perpendicular blast wave geometry on the Imperial Glass Laser it was demonstrated that the clusters are disassembled into a cold gas ahead of the blast wave. This is caused by the multi-keV ions emitted from the initial laser cluster ionisation process. This ion wave propagates outwards and is slowed via the stopping power of clusters. The wave was observed through the lack of formation of a second perpendicular blast wave caused by an abrupt change in the laser absorption between monatomic and cluster gas. The nanoplasma model of cluster explosion dynamics supported the initial ion velocity measurement and the SRIM model agreed with the observed deceleration of the ion driven cluster disassociation wave. SRIM slightly underestimated this slowing, suggesting that clusters may have a higher stopping power than a typical gas. Preliminary experiments using a second laser to preheat the upstream, un-clustered, gas showed an increase in blast wave speed and electron density. Expansion into a pre-heated medium is often seen in supernova expansions into the interstellar medium. A first attempt to resolve the transient electromagnetic fields believed to occur in the blast wave front using proton probing was attempted, however the laser energy on the laser used produced low energy protons which were scattered by the gas

jet. Future work with lower Z gases or higher energy proton beams should overcome this problem however. Finally, the chapter reported on high energy scaling of the blast waves on the Vulcan (Target Area West) Laser. Here highly efficient absorption was observed in the ~ 100 J range. Also effects such as defocusing, saturation and supercritical shocks were observed.

Chapter 6, fully characterised the proton beams generated from laser-foil interactions in the previous chapters to ensure their viability for utilising them as a diagnostic tool. For the first time the low energy (70 keV - 5 MeV) proton spectra was accurately measured for various laser drive energies. The low energy peak was observed not to move from ~ 100 keV as the incident energy increased. Measurements of virtual source size, energy dependent divergence and low energy emittance were also described. The results show that the emission was very close to laminar and for a given energy of protons can be described as being emitted from a virtual source in front of the generating foil, with a half angle linearly dependent on the proton energy. The second half of the chapter described methods for using these proton beams to determine the stopping powers of transient hot and dense states. It was found that the most promising solution was to generate quasi mono-energetic beams from the short pulse laser generated spectrum to enable direct energy loss mapping. Simultaneous time resolved measurements of the samples temperature and density states are also necessary.

We conclude this thesis by discussing some of the future extensions planned to develop the work presented.

7.1 Advanced Optical Probing

The addition of ultra-violet, fourth harmonic probing brings two key advantages over that of green, second harmonic probing. Interferometry, as discussed in §3.3.3, relies on the varying density of the hot matter to cause refractive index variations which in turn causes a fringe shift. However, as the laser intensities increase so often do the electron density gradients. These increased variations in refractive index can induce refraction. This refraction produces phase error in the interferometric analysis and in the extreme case a ray can be refracted away from the collection lens, which will cause a dark region in the image. This ray-bending is analogous to a high emissivity in particle beams as discussed in the previous chapter.

The dark region is caused as the electron density approaches that of a critical density

surface. Sweeney *et al.* determine through ray-tracing that electron densities greater than $n_{critical}/10$ will produce significant refraction effects [152]. Since the critical density is wavelength dependent by halving the wavelength to the fourth harmonic the critical density will change from $4 \times 10^{20} \text{ cm}^{-3}$ to $1.6 \times 10^{20} \text{ cm}^{-3}$. Hence, as higher laser intensities or in shock collisions that also increase the electron densities achieved, moving to ultra-violet interferometry will reduce the error.

A rigorous account of the error induced from ray-bending as a function of the lens speed can be found in work by Gumbrell [153] indicating that for our interferometry a much smaller phase error is induced.

The second advantage comes through the combination of the green and ultra-violet probes - two colour probing. Using second and forth harmonic generation crystals a two colour co-propagating beam can probe a sample from the same angle on a single shot and then be separated using dichroic mirrors before focussing onto a camera. This simultaneous probing can be used to resolve the ray bending [153]. However, if this effect is reduced through the use of a suitable achromatic lens and a low electron density the two colour interferometry can be used to simultaneously measure electron and neutral densities. Throughout this thesis it has been assumed that the bulk of refractive index contribution arises from free electron densities. Recent work has suggested that neutral mass and bound electrons in multiply charged ions can also have a significant contribution to this - for instance at high densities at the shock front. Indeed, the blast wave experiments on the Vulcan laser saw a negative fringe shift, indicating a refractive index greater than one, in several shots as seen in Fig. 7.1. The work by Symes and Rodriguez using the RAPCAL code [75] [2] to show that line radiation dominates Bremsstrahlung by more than three orders of magnitude under our blast wave conditions suggests that bound electrons do have a significant role to play and are likely to be the cause of this fringe shift - significantly affecting the refractive index and interferometric analysis.

The two wavelengths will produce different fringe shifts, ϕ_G and ϕ_{UV} which are a sum of the fringe shifts due to free electrons, ϕ_e , and from bound electrons in neutrals or ions, ϕ_n , where [154] [155]:

$$\phi_e = -r_e \lambda_s n_e L \quad (7.1a)$$

$$\phi_n = \frac{2\pi}{\lambda_s} \frac{\beta}{n_{atom}} n_n L \quad (7.1b)$$

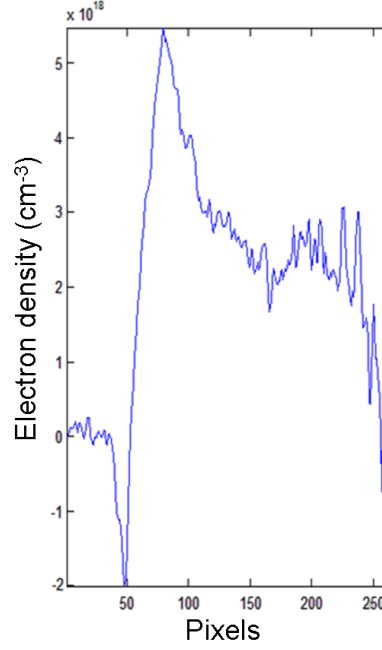


Fig. 7.1: Standard green 527 nm interferometric analysis of a 10 J shot into krypton clusters held at 35 bar backing pressure on the Vulcan laser. The negative electron density feature at 50 pixels as the shock moves from right to left is caused by a fringe reversal which we attribute to bound electrons.

with $r_e = e^2/(4\pi\epsilon_0 m_e c^2)$ the classical electron radius and L the plasma length seen by the probe beam. β is the Gladstone-Dale constant which is wavelength independent and specified for a given neutral species density, n_{atom} , such that the refractive index is $\eta = \beta + 1$. Combining these equations at two wavelengths gives [154] [155]:

$$n_e = \frac{1}{r_e L} \frac{\phi_G \lambda_G - \phi_{UV} \lambda_{UV}}{\lambda_{UV}^2 - \lambda_G^2} \quad (7.2a)$$

$$n_n = \frac{n_{atom}}{2\pi\beta L} \frac{\phi_G \lambda_G - \phi_{UV} \lambda_{UV}}{\lambda_{UV}/\lambda_G - \lambda_G/\lambda_{UV}}. \quad (7.2b)$$

It should be noted that higher wavelengths, for example an infra-red probe, are more sensitive to free electrons whereas the visible wavelengths are more sensitive to bound electrons.

The two equations above allow spatial and temporal recovery of the electron and ion density. A key measurement which is necessary along with this is similarly spatially and temporally resolved electron and ion temperature. A method for obtaining this which time constraints prevented the implementation of in the work reported here is streaked Thomson scattering. Thomson scattering can also provide a second corroborating measurement of

density. Here the narrow band spectrum of incident light from a several Joule pulsed nanosecond laser which is focused onto the sample is compared with the spectrum of the scattered light. The electron temperature can be recovered from the Doppler broadening of the scattered spectrum. Use of two ion species allows recovery of ion temperature. More detailed information can be found in [156] [157].

7.2 Annular Shocks for Jet Production

Continuing work on colliding shocks we can focus on a range of different geometries. One of the key features observed in Fig. 5.5b and the work by James Lazarus using a Fresnel bi-prism to collide shocks is a density enhancement at the collision point [27]. Extending this to the radially symmetric case of an annular shock where the inner surface expands and eventually collides on the laser axis has been modelled by Alberto Marocchino using Gorgon 3D hydrodynamic plasma code, and preliminary experimental work carried out at Imperial by the author.

Using initial conditions of a hollow gas cone heated to 2.5 keV the Gorgon code shows that the resulting expansion into hydrogen drives a cylindrically symmetric Mach stem growth on the laser axis clearly evident after 15 ns as seen in Fig. 7.2. Marocchino suggest that a collision angle of 45° is required to drive the Mach stem. The Mach stem experiment if scaled to higher energy laser systems may allow the generation of high Mach number jets on a relatively small scale laser.

In order to produce the annular focus required to drive the shock wave an *axicon* can be used in combinations with a lens. Here, the radius of the ring at focus is given by $r_{ring} = f \tan(\theta_{out} - \theta_{axicon})$ where θ_{axicon} is the axicon angle which when combined with a lens of focal length f produces a refraction of θ_{out} from Snell's law [158].

Preliminary experiments were carried out using this axicon geometry with room temperature argon gas at 30 bar to produce the clusters on the Imperial College Glass Laser at 700 mJ. Fig. 7.3 shows that the necessary 45° angle of collision was achieved and an electron density enhancement can clearly be seen at the centre of the converging annular shock. However, the Mach stem seen in simulations were not observed. We believe this is due to a need for a higher energy laser - a gas with a higher Z than the simulated hydrogen gas will require higher laser energy to drive the Mach stem. Indeed, during the experiment the blast wave lost definition beyond 40 ns where the Mach stem has not yet

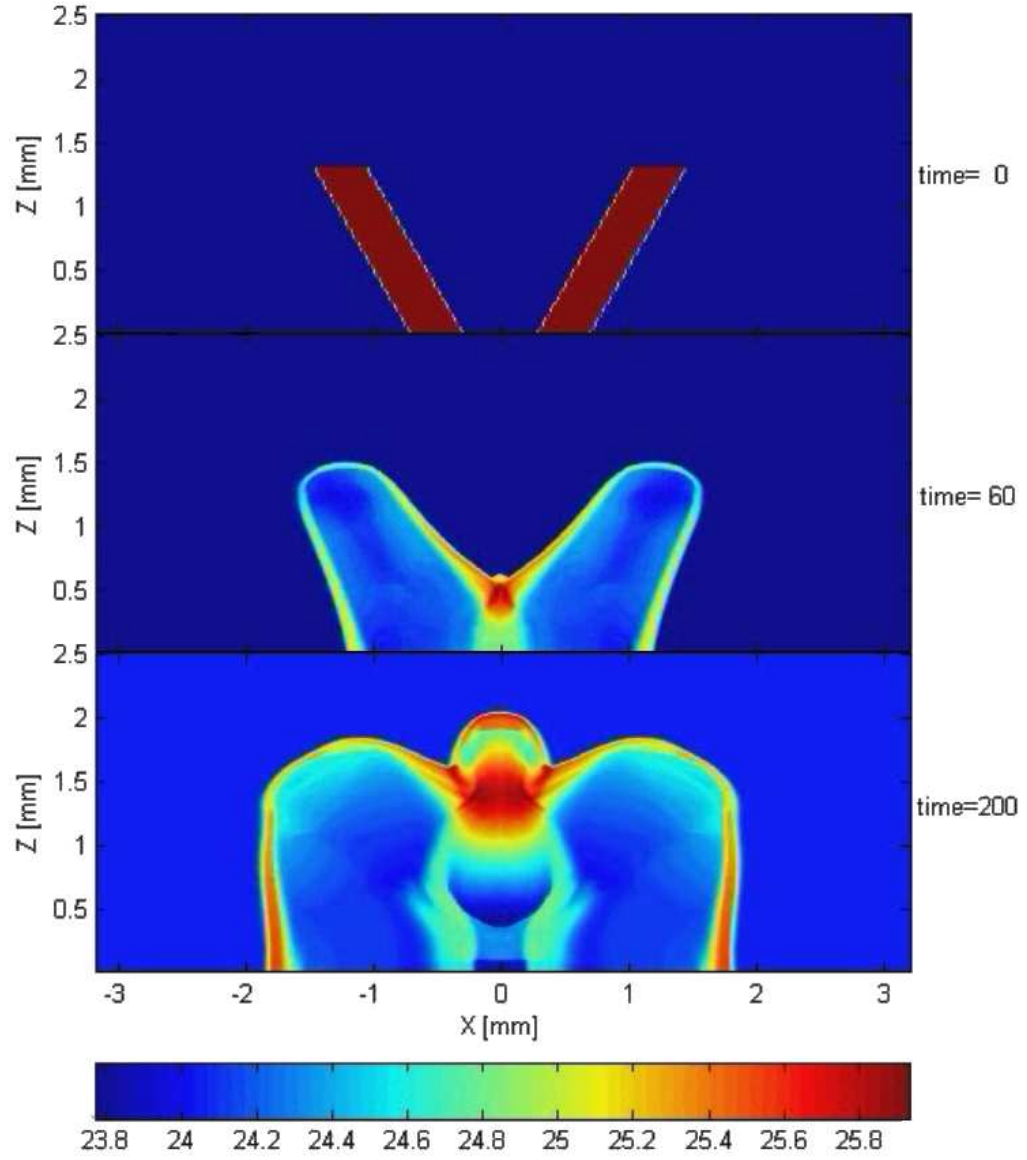


Fig. 7.2: Hydrodynamic simulation of a conical geometry of gas heated to 2.5 keV shown at 0 ns and left to expand. At 60 ns the formation of a rotationally symmetric Mach stem is seen at $x = 0$. This Mach stem then propagates into the surrounding hydrogen gas ($n=10^{10} \text{ cm}^{-3}$) ahead of the converging shock as seen at 200 ns. The colour chart has units of $\log_{10}(\text{electron density (m}^{-3}\text{)})$.

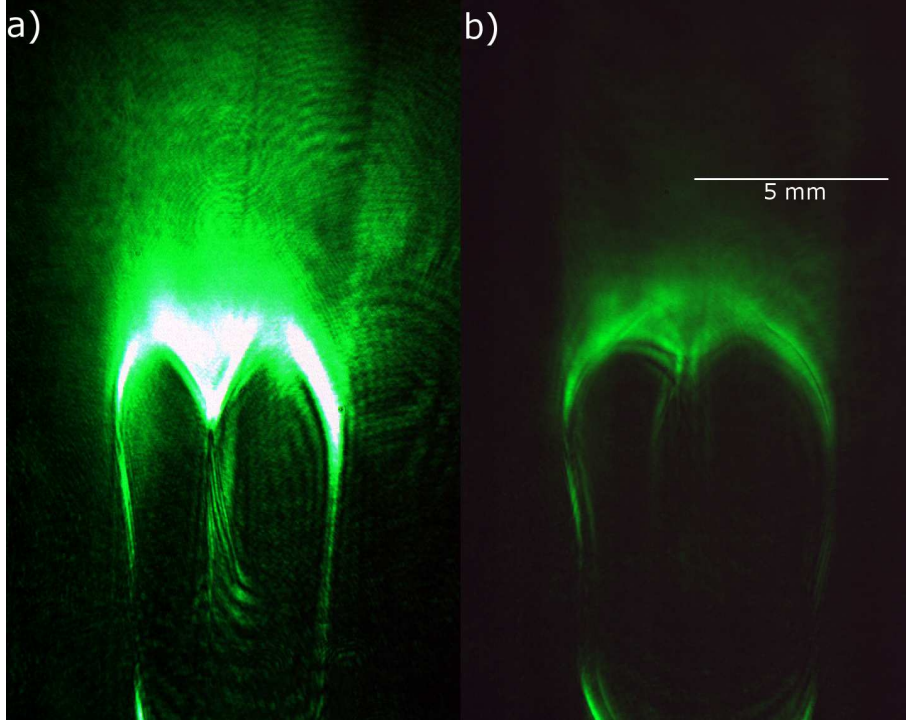


Fig. 7.3: Dark-field schlieren images from preliminary annular blast wave experiment. A 700 mJ, 1053 nm, 0.8 ps laser pulse was focused into argon clusters (30 bar backing pressure) with an axicon to produce a ring focus with a radius of 5 mm. Images at a) 14 ns and b) 21 ns after the laser pulse are shown. The bright regions correspond to high electron density. No evidence of a Mach stem can be seen.

formed in the simulations. Similar experiments in hydrogen clusters and at higher laser drive energies are planned for the future.

7.3 New Laser Capability: Cerberus Laser

Currently the Imperial College Nd:Glass CPA Laser is being replaced by the Cerberus Laser OPCPA system. This new multi-beam line system will bring new capabilities and higher energies. This will allow experiments such as blast wave proton probing, annular blast wave production, development of a laser driven quasi mono-energetic ion beam and further experiments discussed in this thesis to be carried out on a university based laboratory system with a relatively high repetition rate. The Cerberus system has also been designed to have access to the MAGPIE chamber such that high energy laser probing of the plasmas produced by this pulsed power generator can be carried out - particularly laser driven proton probing, x-ray imaging and a visible Thomson scattering probe.

Cerberus, after the multi-headed Greek mythological hound which guards the gates of

the Underworld to prevent escape across the river Styx, is so named due to the multiple beam lines produced from one front-end. Specifically the Cerberus laser has one short pulse >20 J, ~ 500 fs beam line and three ‘long’ pulse >50 J, nanosecond beam lines. This is a unique laser capability outside of national facilities, particularly in combination with MAGPIE, and is supported by components acquired from the Central Laser Facility and the HELEN laser. Fig. 7.4 shows the current schematic of the Cerberus laser. As can be seen the two key upgrades from the Imperial College CPA Nd:Glass are an OPCPA scheme and several higher energy, large aperture amplifiers.

7.4 Concluding Remarks

Laser driven ion beams are a unique and versatile tool due to their relatively high penetration, short pulse length and electromagnetic field susceptibility. The experimental and numerical investigations presented in this thesis show that these ions provide a means to access the warm dense matter regime with some confidence. Furthermore, they allow determination of the ion stopping power in both this transitory regime and the high energy density plasma regime. The high laser absorption of cluster gases allows access to the high energy density regime through the production of radiative, and non-radiative, blast waves. This thesis has further characterised these blast waves and their propagation dynamics with novel geometries and identified proton probing as an important tool for future investigations of their self generated field structures, and their ability to entrain and compress external fields. It is believed the work described here will provide fundamentally important insight into key physical processes including the heating process on the journey to ‘ignition’ for laser fusion targets, and also astrophysical phenomena scaled to the laboratory.

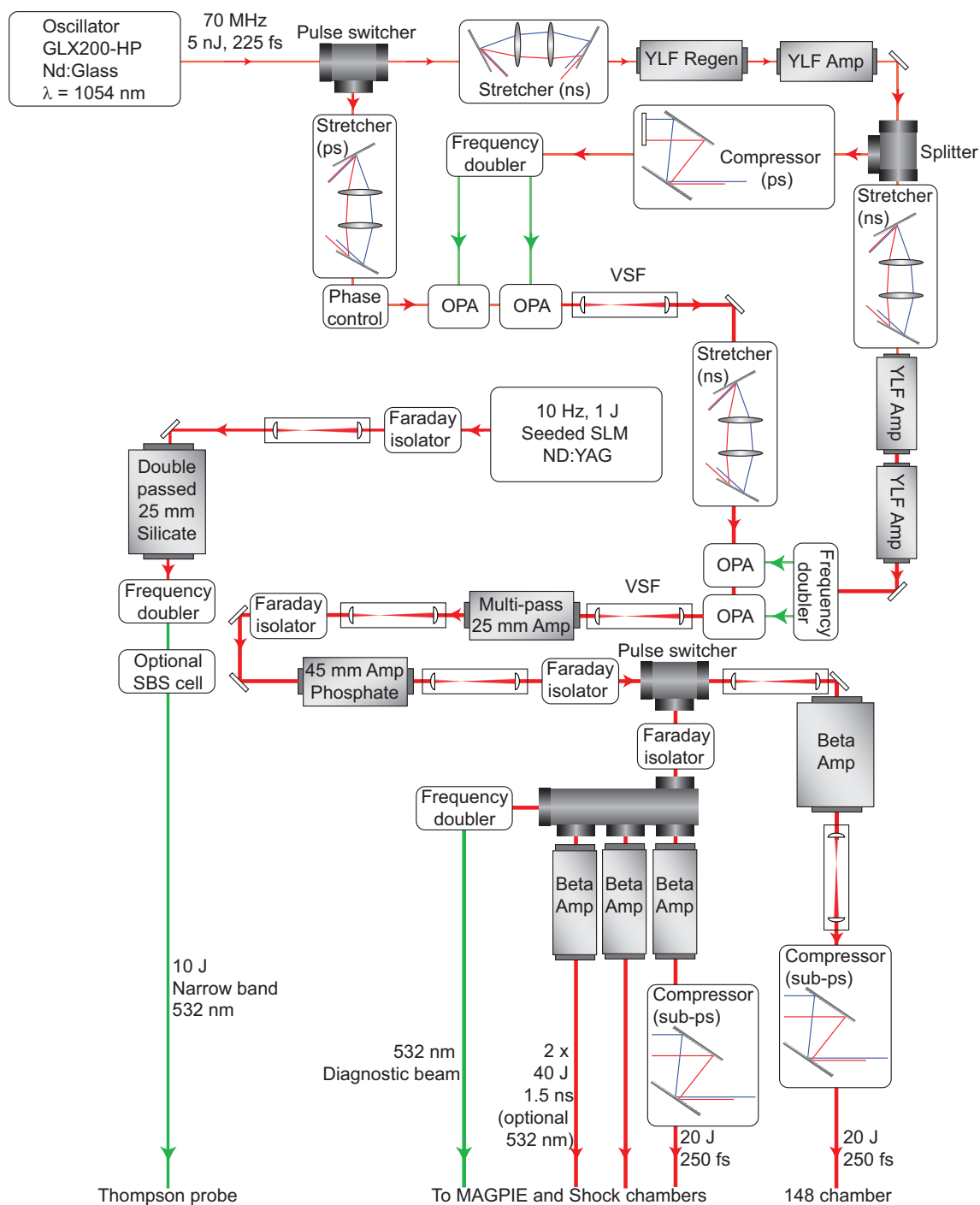


Fig. 7.4: Schematic of Cerberus multi-beam laser system under development by R. A. Smith and co-workers at Imperial College.

Appendix A

Derivation of Magnetic Spectrometer Error

As a non-relativistic ion propagates through the entrance slits of the magnetic spectrometer we can ignore the non-uniformity and fringe effects of the magnetic fields for the error calculations. Hence, the problem reduces to one of a charged particle in a uniform magnetic field where the equation of motion is:

$$m\dot{\boldsymbol{\nu}} = q(\boldsymbol{\nu} \times \mathbf{B}) \quad (\text{A.1a})$$

$$\dot{\boldsymbol{\nu}} = \omega_L \times \boldsymbol{\nu} \quad (\text{A.1b})$$

where $\omega_L = \frac{|q|\mathbf{B}}{m}$ is the Larmor, or gyro, frequency. The force on the particle is perpendicular to the velocity of the particle - therefore the energy remains constant and any velocity in the field direction (z -axis) is unchanged.

The perpendicular speed is thus:

$$\nu_{\perp} = \sqrt{\nu^2 - \nu_{\parallel}^2} \quad (\text{A.2a})$$

$$\boldsymbol{\nu}_{\perp} = \nu_x \hat{\mathbf{x}} + \nu_y \hat{\mathbf{y}}. \quad (\text{A.2b})$$

Using Eq. (A.1b) we then obtain the familiar relations [159]:

$$\dot{\nu}_x = \omega_L \nu_y \quad \dot{\nu}_y = -\omega_L \nu_x \quad (\text{A.3a})$$

$$\ddot{\nu}_x = -\omega_L^2 \nu_x \quad \ddot{\nu}_y = -\omega_L^2 \nu_y \quad (\text{A.3b})$$

which we recognise as simple harmonic motion with a phase offset between the x and y

axis to produce circular motion with the solution:

$$\nu_x = \nu_{\perp} \sin(\omega_L t + \phi) \quad (\text{A.4a})$$

$$\nu_y = \nu_{\perp} \cos(\omega_L t + \phi) \quad (\text{A.4b})$$

where $\phi = \arctan(\nu_x^0/\nu_y^0)$ and from the initial conditions of the entrance slit we can see that $\phi = \frac{\pi}{2}$. Further integration returns the solution to the equation of motion at a given position:

$$x = x_c - r_L \cos(\omega_L t + \phi) \quad (\text{A.5a})$$

$$y = y_c + r_L \sin(\omega_L t + \phi) \quad (\text{A.5b})$$

where these equations give the relation $(x - x_c)^2 + (y - y_c)^2 = r_L^2$. The variable $r_L = \frac{\nu_{\perp}}{\omega_L} = \frac{m\nu_{\perp}}{qB}$ is known as the *Larmor radius* (or gyroradius). x_c and y_c define the coordinates of the centre of the circle with radius r_L , known as *the guiding centre*. Using the initial boundary conditions we can set $x_c = x_0$ and $y_c = y_0 - \frac{\nu_{\perp}}{\omega_L}$ where $\mathbf{r} = x_0\hat{\mathbf{x}} + y_0\hat{\mathbf{y}}$ is the initial ion position at the entrance slit.

Hence, for a given position on the detector plane (x, y) we can find the ion energy, E :

$$\begin{aligned} E(x, y) &= \frac{1}{2}m\nu_{\perp}^2 \\ &= \frac{1}{2}m\omega_L^2[(x - x_c)^2 + (y - y_c)^2] \\ &= \frac{|q|^2 B^2}{2m}[(x - x_0)^2 + (y - (y_0 - \frac{\nu_{\perp}}{\omega_L}))^2] \\ &= \frac{|q|^2 B^2}{8m}[(x - x_0)^2 + (y - y_0)^2]^2 / (y - y_0)^2 \end{aligned} \quad (\text{A.6})$$

Sources of error will arise from the ion beam divergence and the finite slit width limiting the resolution. The divergence is minimised by using double slits and placing the source at a distance from the spectrometer so that the beam is well collimated. The slit must have a finite width, Δy , which induces a spread of energies, δE , to be detected at one point on the detector plane where:

$$\delta E = E(x, (y + \frac{\Delta y}{2})) - E(x, (y - \frac{\Delta y}{2})). \quad (\text{A.7})$$

Using the dimensions of the spectrometer we can plot the normalised error, $\delta E/E$, along the horizontal and up the vertical detector plane as seen in Fig. A.1 for the typical 250

μm slit size. The error profile has three distinct regions. At ~ 1650 keV the turn from the horizontal plane to the vertical detector plane can clearly be seen. At ~ 100 keV and below the dispersion becomes anomalous as the ion orbit begins to bend back on itself. In other words the orbit has gone through more than 90° .

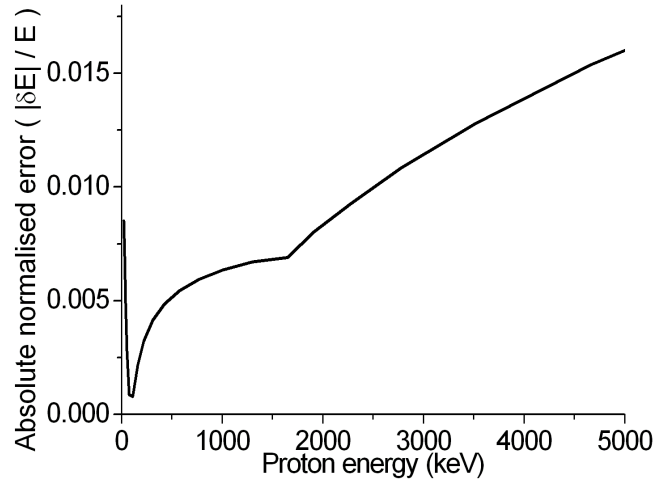


Fig. A.1: Normalised error of the magnetic spectrometer arising from a finite slit width. A standard slit width of $250 \mu\text{m}$ is used. The profile has three distinct regions. The right region shows the error for the vertical detector plane. The middle region shows the error for the horizontal plane. The far left region shows the beginning of the anomalous error region where the ion orbit passes through 90° and the dispersion rapidly increases.

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