Channeling Experiments on OMEGA EP

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Biographical Sketch

The author was born in Amesbury, Massachusetts in 1985. He attended the University of Rochester supported by a Bausch and Lomb Science scholarship and the Kraft Foods science and engineering scholarship. He graduated with high distinction in Engineering in 2008. He then matriculated into the doctoral program and began his studies at the Laboratory for Laser Energetics. In 2010, he received a Master of Science degree in Mechanical Engineering and remained at the University of Rochester to fulfill the requirements of the Doctoral degree in Mechanical Engineering. His thesis work was performed at the Laboratory for Laser Energetics under the direction of Prof. D.D. Meyerhofer and Dr. W. Theobald.
Publications

Publications and selected professional conference presentations of the author during his time at the laboratory include:


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Abstract

The creation of stable, straight channels in dense plasma with high intensity laser beams is of interest to a number of applications in laser-produced plasma science. This includes the generation of MeV to GeV electron beams in laser wakefield accelerators (LWFA)\(^1\), coherent and brilliant X-ray generation from electron betatron radiation\(^2\), and fast ignition of inertial confinement fusion targets (FI ICF)\(^3\). This thesis studies the creation of channels through dense preformed plasma on the OMEGA EP laser system\(^4\). A new optical probing method was developed to image the channeling process. The experiments show that the channel bores faster than the ponderomotive hole boring velocity of Wilks \textit{et al.}\(^5\) predicted by a balance of light pressure at the channel front. In the wake of the channeling pulse, a shock-driven parabolic density profile exists in the radial direction, allowing for the guiding of a subsequent high intensity pulse to an embedded steep gradient. A structure of this type is beneficial to the applications mentioned above. The results of this study can be extended to an integrated FI experiment on OMEGA, planned for the upcoming year.

Current laser systems are able to create high energy density (HED) systems where the radiation-pressure in the laser field totally overwhelms the thermal pressure of the material. In this radiation-pressure dominated regime, the radiation-pressure, \(2I/c\), where \(I\) is the intensity of the radiation field and \(c\) the speed of light, is orders of magnitude greater than the thermal pressure, \(nk_BT\), where \(n\) is the material density and \(k_BT\) is the specific energy of the system. This radiation-pressure driven flow modifies
the plasma density in the region of the laser. This work presents measurements of a channel or density cavitation bored into an inhomogeneous HED plasma.

A high-energy chirped-pulse amplification (CPA) laser beam ($I \sim 10^{18}$ W/cm$^2$) was used to drive a radiation-pressure driven blast-wave into the plasma. A 10-ps, 10-mJ, 3.6 $\mu$m resolution optical probing system was developed to image the interaction with unparalleled resolution with high spatial and temporal resolution. In a second experiment, a second, co-propagating high-intensity pulse was guided through the blast-wave prepared channel established by the first pulse. The pulse is transmitted efficiently and maintains a spatial profile close to the vacuum mode. This is most likely due to the intense heating experienced by the plasma during the channeling beam.

The experimental data show a self-similar cylindrical blast-wave developing from the channeling laser. The forward velocity of the channel is measured and is quantitatively consistent with a radiation-pressure driven flow model that predicts the propagation velocity. The residual plasma in the channel is heated to multi-MeV temperatures (11,000,000,000 K) and the long time expansion ($\sim 500$ ps) is consistent with 2-D radiation-hydrodynamics simulations.

This work shows the creation of radiation-pressure driven flows in plasma. The creation of these conditions may be scaled to a regime where the dimensionless ratio of radiation pressure to thermal pressure, $2I/nTc$ is $>> 1$. Conditions in this regime are found in physical conditions ranging from stellar interiors and radiation generated winds, to the formations of “photon bubbles” in very hot stars and accretion disks$^{6,7}$ and in high-energy astrophysical environments$^8$. 
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Chapter 1

Introduction

The creation of high energy density physics (HEDP) in the laboratory was facilitated by the invention of the laser. The interaction of high-power lasers with matter creates material with high temperature and density simultaneously. In a tiny fraction of a second, laser systems can create material at temperatures and pressures that prevail inside astrophysical objects throughout the universe. This chapter motivates the thesis, through an introduction to HEDP and laser driven nuclear fusion. The chapter concludes with a summary of the work presented in this thesis.

1.1 High energy density physics

The nascent field of HEDP was ushered in with the dawn of the thermonuclear age. HEDP is separated from traditional condensed matter physics and plasma physics by the energy density of the system. The energy density associated with HEDP is $10^{11}$ J/cm$^3$ or an equivalent pressure of 100 GPa or greater. High power transforms matter into HEDP. It may provided by be electrical, laser, or thermonuclear power. HEDP spans a vast scale in time and space, providing insights into the development of galactic magnetic field, the interior structure of stars and giant planets, the seeding
of lighting bolts on earth\textsuperscript{23,24}, creation of pair plasmas in the laboratory\textsuperscript{25}, and the development of new particle accelerators\textsuperscript{26,27,28}, and the creation of conditions capable of nuclear fusion in the laboratory\textsuperscript{29}. The last forty years have seen unprecedented developments in the technology enabling the observation and measurement of HEDP phenomena.

1.2 Nuclear Fusion

A major effort of HEDP science is devoted to the understanding of nuclear fusion in the laboratory and the ability to drive fusion reactions in a controlled manner. Nuclear fusion is a nucleon synthesis process where two chemical elements assemble a lighter nucleus from their respective nucleons. The nucleus of an atom is held together by the strong force between protons and neutrons. The result of this force is a binding energy representing the mechanical work required to separate the nucleus into it’s constituents. The isotopes of the lightest element, hydrogen, have a lower binding energy than it’s higher mass neighbor on the periodic table, helium. Heavy elements such as uranium and plutonium possess a lower binding energy than elements in the center of the periodic table such as iron or nickel. The process of a heavy nuclei disassembling into products with lower atomic numbers is called nuclear fission, while the process when light elements form a higher atomic number element is called nuclear fusion. In both cases, energy is released either by fusion or fission when the nucleons are rearranged to a nuclear configuration with a higher binding energy per nucleon and hence lower mass. Figure 1.1 shows the example of the fusion of a deuteron and triton to a helium nucleus and a neutron.

The electrostatic repulsion of the Coulomb force between two nuclei must be overcome for the nuclei to fuse. The attractive nuclear force is much stronger than the
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Coulomb force, but drops off rapidly with distance. For the nuclei to overcome the Coulomb barrier, enough kinetic or thermal energy to overwhelm the repulsion must be supplied so that the colliding atoms can fuse. The thermal energy of atoms is measured as temperature in a macroscopic sense. The temperature required for fusion to occur is high, in excess of $10^7$ K. Thermonuclear fusion plasmas are difficult to contain in the laboratory. Maintaining the fuel at thermonuclear temperatures is a challenge because it is far above the melting point of any container that can be built. This is one of the reasons why a fusion power plant has so far not been realized.

Fusion reactions occur naturally at the center of all stars. Stars are brilliant celestial spheres of plasma held together by their own gravity. The source of plasma confinement is the gravitational force that holds the star together. As the star contracts under the force of gravity, the hydrostatic pressure of the star increases. The star contracts until the internal pressure supports the star from further gravitational contraction.

The energy that is necessary to overcome the Coulomb force of the two nuclei increases with the atomic number of the reactants. The nuclei with the lowest barrier to overcome for nuclear fusion are the isotopes of hydrogen: ordinary hydrogen ($^1H$ or
p), deuterium ($^1_2H$ or D) and tritium ($^3_1H$ or T). Nuclear reactions of interest with these elements are

$$D + D \rightarrow T + p + 4.0 \text{ MeV},$$  \hspace{1cm} (1.1)

$$D + D \rightarrow ^3\text{He} + n + 3.3 \text{ MeV},$$  \hspace{1cm} (1.2)

$$D + T \rightarrow \alpha + n + 17.6 \text{ MeV}. \hspace{1cm} (1.3)$$

The likelihood of the reactions occurring is a function of temperature and density of the interacting nuclei. The probability of reaction of a nucleus per unit path length inside a target is given by $n\sigma$, where $n$ is the target density and $\sigma$ is the fusion cross-section. The cross-section increases for D-T reactions with increasing temperatures up to 100 keV.

High temperature is one requirement for efficient fusion of light elements. For the net generation of energy, one must consider the energy input required to maintain the reactants at the necessary temperature $T$, the required particle number density $n$, and the time the reactions take place $\tau$. The fusion reactions must release more energy than is required to heat the plasma and account for the unavoidable losses due to radiation. Lawson$^{30}$ showed that these parameters form an ignition condition that gives the threshold for net energy production from fusion reactions. This condition is most commonly written as

$$nT \tau > 5 \text{ atm s}, \hspace{1cm} (1.4)$$

where $nT$ is the pressure of the fuel, measured in atmospheres and $\tau$ is the confinement time in seconds. This criterion forms the basis for break-even energy production from
fusion reactions. The goal of laboratory thermonuclear fusion experiments is to create a plasma that meets this requirement. The challenge is to satisfy both the pressure and time constraints simultaneously\textsuperscript{29,31}. In the most familiar self-sustaining nuclear fusion reaction, the sun has a confinement time measured in billions of years rather than seconds. The pressure and density are orders of magnitude higher than the highest pressures and densities achievable in the laboratory, leading to a twenty-five order of magnitude excess over the breakeven ignition condition for the conditions deep in the Sun ($T = 15 \times 10^6$ K, $n = 1 \times 10^{32}$ m$^{-3}$, and $\tau = 3 \times 10^{14}$ s)\textsuperscript{32}.

There are two approaches to contain the thermonuclear fuel in the laboratory, magnetic and inertial confinement. In magnetic confinement, the plasma is constrained to travel along imposed magnetic field lines. The magnetic field forces the ions and electrons along the field lines, which keeps the hot fuel away from the walls of a containment vessel. In inertial confinement fusion, the confinement time is determined only by the inertia of the compressed fuel without any additional effort to contain the fuel. The plasma is only hot and dense simultaneously for the time it takes for a sound wave to travel across the fuel volume. All of the thermonuclear energy is released in a short burst. The inertial method is discussed further in the next section.

### 1.3 Inertial Confinement Fusion

The process of Inertial Confinement Fusion (ICF)\textsuperscript{33} consists of four parts to assemble a mass of fusion fuel to the temperature and density required by Lawson’s criterion. The four phases are acceleration, coasting, deceleration, and stagnation. The capsule of DT-ice is driven inward by reaction forces from the momentum transfer of the outer ablator burning off the capsule surface. In the first two images of Figure 1.2, the rocket effect accelerates the shell inward with a velocity of $>3 \times 10^7$ cm/s. The
capsule implodes to a radius that is 20-30 times smaller than the initial capsule radius. The volumetric compression ratio of the capsule is $V_i/V_f \sim 40000$. After about half of the implosion time, the driver turns off which initiates the second [Fig. 1.2(b)] phase of the implosion. During this time the shell flies inward building up pressure inside of the shell. This phase occurs between the second [Fig. 1.2(b)] and third [Fig. 1.2(c)] image in Figure 1.2. At a point during the coasting, the pressure in the center reduces the momentum of the inward flying shell. During the next step, deceleration, the inward flying shell begins to slow as the kinetic energy of the shell is converted to internal energy of the heated gas in the hot spot formed in the center of the shell. At this point the pressure is the highest during the implosion, which is called stagnation because of the shell’s momentary pause before disassembling. It is at this moment that the compression of the shell reaches a maximum. The shell’s trajectory reaches a turning point from imploding to exploding. Profiles of the temperature and pressure of the compressed pellet are sketched in Figure 1.2(d). The fuel assembly has a hot center of gas that is heated by the $PdV$ work of the imploding shell. At stagnation, the temperature in the hot core is high enough so that D and T particles begin to fuse releasing neutrons and $\alpha$-particles. The $\alpha$-particles are stopped in the hot spot because sufficient areal density is provided for the charged particles that undergo many Coulomb collisions. The areal density refers to the product of $\rho R$ where $\rho$ is the mass density of the plasma and $R$ is the average radius of the confined plasma. The areal density is a proxy for the $\alpha$-particle stopping power of the plasma. The $\alpha$-particles transfer their kinetic energy to the hot spot, heating the fuel and launching an expanding burn wave that propagates from the center of the capsule outward. This process is limited to the brief duration of hydrodynamic stagnation, which is $\sim 0.1$ ns for small fuel capsules.

The ICF capsule may be driven using overlapped symmetric laser illumination (direct-drive) or by a bath of thermal x-rays generated in a high-Z container (indi-
1.3. INERTIAL CONFINEMENT FUSION

Laser irradiation DT fuel-filled pellet
Expanding coronal plasma
Imploding pellet
Compressed pellet
Temperature
Radius
Mass density
Hot spot
Burn wave
Stagnation / Ignition

Figure 1.2: In the ICF process, the fuel is heated by the PdV work of the rapidly collapsing shell. At the time of peak compression, the hot spot in the center begins to fuse and release alpha particles. The 3.5 MeV alpha particles are stopped by the surrounding dense shell and deposit energy which heats the entire volume to thermonuclear temperatures initiating a chain reaction.

rect drive) by laser irradiation. Figure 1.3 shows a schematic representation of the two capsule drivers. Both methods couple energy to the capsule through ablation of the outer surface of the shell. The direct-drive energy efficiency is greater than indirect drive because coupling of the x-ray drive to the target is low due to losses in heating the container and radiation escaping out of the ends of the container. The targets for direct drive are simpler because they do not require the high-Z container that is needed for indirect drive. Indirect drive provides a more uniform illumination than direct drive because the spatial inhomogeneity of the x-ray bath is less than that of a high-power laser beam. Efforts have been successful in smoothing out the laser non-uniformities.

The ICF method has demonstrated tremendous thermonuclear gains in uncontrolled fusion experiments driven by a nuclear explosion. It is difficult to scale this down to reasonable sizes and drivers that can be used in a laboratory or in a power plant. For thermonuclear ignition of a capsule, the shell must reach a velocity of at least $3 \times 10^7$ cm/s. The shell must be thin for the driver to accelerate the fuel mass to that velocity and maintain a high compression ratio. Thin shells with a high aspect ratio (shell radius divided by the shell thickness) are prone to hydrodynamic instabilities that might prevent ignition. Although significant progress has been made in recent ICF experiments by demonstrating the onset of $\alpha$-heating and reaching fusion output
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Figure 1.3: A schematic representation showing the irradiation configuration for both direct (a) and indirect (b) drive.

exceeding the kinetic energy of the inward flying shell 29, ignition has so far not been achieved. Alternative ignition schemes are fast ignition (FI) 3 and shock ignition 36 that rely on an additional step by first assembling the fuel to a high density with a sub-ignition implosion velocity and then delivering the required energy for ignition from a separate external source. Fast ignition uses an isochoric model (uniform density) that compresses far more mass to a high areal density. By burning the extra mass, the yield and thermonuclear gain are considerably larger than that of the conventional isobaric (hot-spot or uniform pressure) model. The FI scheme uses a beam of charged particles generated by a powerful short-pulse laser to ignite the fuel. For a fixed amount of compressional work, FI might be more efficient in the burning of thermonuclear fuel and reach higher gains, provided the fuel is ignited.

1.4 Fast Ignition

In the original proposal for FI 3, an intense laser pulse creates an evacuated channel in the coronal plasma of an imploded ICF target. The channel provides a path through which a second, co-axial high-intensity laser pulse propagates. The energy of
the second intense pulse is partially converted into high energy electrons at the end of
the channel. These penetrate and heat the imploded fuel assembly at the time of peak
compression, providing the energy to initiate the thermonuclear chain reaction. This
scheme has benefits compared to traditional hot-spot ICF, allowing stable, low-velocity
implosions while achieving high gains\textsuperscript{37}. The increase in potential gain as a function
of drive energy is illustrated in Figure 1.4\textsuperscript{9,10}. The dark blue and light blue regions in the
plot represent the variation of the target gain with the driver energy for direct drive and
indirect drive, respectively. The two light blue bands are simulations for two implosion
velocities. The dashed lines indicate what gain corresponds to what driver energy for
a fixed total fusion yield (10 MJ, 100 MJ, and 1000 MJ). If one uses the driver energy
solely to compress the fuel and provides a separate source of energy for ignition, then
the black curves give the potential gain for these implosions. The required amount of
laser energy to ignite the imploded assembly, the maximum density, and the estimated
laser intensity of the ignitor beam are written on the side of each curve. The amount
of fast heating energy required to ignite the target depends on the fuel’s density. The
curves show the difference between 150, 200, and 300 g/cc final target density, which
requires 330, 200, and 92 kJ of fast ignition laser energy, respectively. One way to
supply the ignition energy into a small volume during the time of peak compression,
which is only tens of picoseconds, is to use powerful chirped pulse amplified lasers
(CPA)\textsuperscript{38,39} that produce a beam of charged particles\textsuperscript{40}.

The advantages of FI over conventional ICF are realized by the differences in cap-
sule implosions. To create the hot-spot necessary for conventional ICF, shell implosion
velocities of at least \(3 \times 10^7\) cm s\(^{-1}\) are required\textsuperscript{41} or there is insufficient kinetic energy
in the shell available to ignite the target. To achieve a high shell velocity for conven-
tional ICF, a high aspect-ratio (ratio of shell radius to thickness), low mass target must
be used. Such a target is more likely to break up in-flight\textsuperscript{42}. In contrast, a FI capsule
can be imploded with a lower velocity (around $1.5 \times 10^7$ cm/s), using a higher mass target with lower aspect-ratio that possesses better hydrodynamic stability than a conventional hot-spot ignition target. More target mass imploded with the same amount of driver energy results in higher gains\textsuperscript{37}, provided the fuel is ignited. Here, the gain only takes the driver energy into account and it is assumed that ignition can be achieved with a reasonable short-pulse laser energy.

Several options of ignition pulse delivery are being studied for use in FI\textsuperscript{43}. The injection of an ignition pulse directly into the corona of a FI target has been investigated. The ignition pulse is affected by a number of instabilities in coronal plasma, limiting the effectiveness of hot electron generation and subsequent core heating\textsuperscript{44}. An alternate scheme has been examined, where the low density plasma is kept from interacting with the ignition pulse by a metal cone cast into the target. This scheme has been demonstrated with remarkable success in experiments\textsuperscript{45}. However, the cone-in-shell scheme has its limitations due to preplasma creation in the cone from pre-pulses of the
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Figure 1.5: Schematic of fast ignition with channeling. In this scheme an intense laser is employed to evacuate a column of plasma in the corona to allow a trailing, ultra-intense laser to be focused at the critical surface. Hole-boring pushes the critical surface closer to the center of the target, where the pulse’s energy is converted to high energy electrons that that stream into the core and ignite the fuel. Image reproduced from[11]

ignition pulse[46], the mixing of cone material into the imploded fuel assembly[47], and the long-term cost effectiveness of producing a complicated target. Prepulses inherent in CPA systems can preionize the cone material and decrease the coupling efficiency of the ignitor pulse beam energy to fast electron energy[48]. Contamination of the hot dense fuel assembly with metals increases the radiated energy from the imploded core in the form of x-rays from bremsstrahlung or line radiation from high-Z materials. These issues are avoided in the channeling concept where one pulse produces an empty plasma channel to guide an ignition pulse as described in the initial FI reference[3]. This has the advantage of unimpeded transmission of the ignitor without the drawbacks of a metal cone.

Recent particle-in-cell simulations predict that the ponderomotive pressure of an intense laser pulse can evacuate a channel in the coronal plasma, eliminating the need for a cone[49]. This channel allows the transmission of a trailing, ignition pulse (typi-
Figure 1.6: High energy laser facilities providing both long and short pulse beams. Prior to the activation of the OMEGA EP laser system, GEKKO-XII\textsuperscript{12} was the only facility to conduct integrated FI experiments. Other FI relevant laser systems will be FIREX and NIF-ARC. Figure reproduced from\textsuperscript{13}.

cally 10-20 ps duration) without any of the losses associated with propagation through an extended tenuous coronal plasma of the implosion. Figure 1.5 shows a schematic representation of the FI with channeling scheme. A channel drilled by lasers has been shown to act as a waveguide for a trailing pulse\textsuperscript{50}. The channel has a low enough residual plasma density so that laser-plasma instabilities are not excited\textsuperscript{44}.

The laser parameters required to reach the critical density \((n_e = \varepsilon_0 m_e \omega_l^2 / e^2\), where \(\varepsilon_0\) is the free space permittivity, \(m_e\) is the electron mass, \(\omega_l\) is the laser angular frequency, and \(e\) is the elementary charge\) in an inhomogeneous FI-relevant plasma is a pulse duration of \(\sim 100\) ps, an energy of \(> 2\) kJ, and an intensity of \(5 \times 10^{18}\) W cm\(^{-2}\).\textsuperscript{49} This parameter regime is covered by the OMEGA EP laser system\textsuperscript{4}.

OMEGA EP\textsuperscript{4} provides a platform to conduct channeling experiments in an FI relevant coronal plasma. Prior to OMEGA EP, integrated FI research was conducted on GEKKO-XII\textsuperscript{51}. Figure 1.6 shows several current and proposed FI laser research systems. The long-pulse energy is given by the x-axis and the short-pulse energy by the
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Figure 1.7: Density profile from a 1-D spherical FI implosion and a 2-D planar foil plume that shows the similarity in scale length between the two. The coronal plasma density profile is similar to the expected profile in a FI target with a long density scale length of \(L_s = 430 \, \mu \text{m}\) at the critical density of IR light.

y-axis. OMEGA EP operates at higher short and long pulse energies than GEKKO-XII\(^{12}\) (GXII) and will probe unexplored territory for FI. Two other laser systems will operate in a similar regime as OMEGA EP, FIREX-1\(^5\), an upgrade to the GXII laser, and NIF-ARC\(^5\), a planned short-pulse capability for the National Ignition Facility. The ‘next-generation’ laser systems that are capable of performing FI experiments at relevant plasma conditions are shown in green in Figure 1.6. Figure 1.7 shows simulated density profile for a FI implosion calculated by the 1-D code LILAC\(^5\) and an on-axis density profile from a planar foil irradiated with 4 kJ of UV light from OMEGA EP. The plasma from the planar foil was simulated with the 2-D radiation-hydrodynamics code DRACO\(^5\). The density and its scale length from the planar foil (\(L_s^{-1} = \nabla n_e/n_e\)) are similar to what might be found in an implosion target at the time of peak compression\(^5\).

Channeling provides an opportunity to address challenging physics issues. High
order phase modulations in the laser beam\textsuperscript{57} or induced by the coronal plasma\textsuperscript{58} in the channeling beam can seed filaments that may merge into a single channel or spray the energy into a diffuse cone. In addition to filamentation, the low-density channeling is subject to geometric instabilities that bend or kink the channel into a non-straight path, which might divert the beam from the ignition region\textsuperscript{59}.

In summary, FI promises higher gains and more stable implosions than conventional ICF targets. The relaxation of complexity in fuel compression comes at the expense of the relatively unknown physics of energy transport of the fast electrons through the coronal plasma. Channeling FI allows for symmetric implosions that fully exploit to benefits of low velocity, high density implosions. Better understanding of the physics of channeling in hot, long scale length plasmas might lead to a baseline target design for advanced ignition ICF targets.

1.5 Previous Experiments

Recently, channeling experiments have been conducted on a number of laser systems\textsuperscript{60,61,62}. Channels were observed using optical techniques\textsuperscript{62}, x-ray radiography\textsuperscript{61}, and high energy proton probing\textsuperscript{63,64,65}. The prior experiments have typically been outside of the scope of FI relevant plasmas in either channeling pulse duration (too short)\textsuperscript{60}, plasma extent (too small)\textsuperscript{61}, plasma density (too low)\textsuperscript{65}, or density scale length (too homogeneous)\textsuperscript{64}. Experimental studies of channeling in high-density plasma with ultra-intense lasers have been carried out on several laser systems from around the world\textsuperscript{1,60,61,63,64,65,66}. These experiments are summarized in Table 1.1. Because of each laser facility’s unique architecture and instrumentation, experiments at each facility have probed their own set of experimental conditions, specifically laser energies, pulse durations, plasma density, and plasma scalelength. For the previous
channeling studies, not all of the parameters are those anticipated on a FI implosion. This section summarizes the major experimental studies in FI relevant channeling to date and outlines where the research presented in this work builds on understanding from previous experiments.

Early work by Young\textsuperscript{66} found that the channeling duration of the laser pulse was important to the propagation through the plasma. The plasma was generated by exploding a thin (0.5 μm) CH foil with a 50 J, 1 ns square laser pulse, leading to a peak density of $2 \times 10^{20}$ cm$^{-3}$. They varied the channeling pulse duration from 100 to 1000 ps in a series of experiments, and found increased laser transmission through the plasma for the longer pulses. The laser intensity of the channeling pulse was held constant at $5 \times 10^{16}$ W/cm$^2$ and had a wavelength of 532 nm. They used an interferometer, set up orthogonal to the channel axis, to measure plasma density perturbations. A collection system was used to gather light transmitted through the plasma and calculated the

<table>
<thead>
<tr>
<th>System</th>
<th>Ref.</th>
<th>$E_L$ (J)</th>
<th>$T_L$ (ps)</th>
<th>$n_e$ (cm$^{-3}$)</th>
<th>Diagnostics</th>
</tr>
</thead>
<tbody>
<tr>
<td>JANUS</td>
<td>\textsuperscript{66}</td>
<td>50</td>
<td>100 - 500</td>
<td>$2 \times 10^{20}$</td>
<td>Optical Interferometry, Transmitted Energy Calorimeter</td>
</tr>
<tr>
<td>GXII</td>
<td>\textsuperscript{61}</td>
<td>140</td>
<td>100</td>
<td>$10^{20}$</td>
<td>X-ray Grid Image Refractometry, X-ray Pinhole Imaging</td>
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<tr>
<td>VULCAN</td>
<td>\textsuperscript{60}</td>
<td>10</td>
<td>1</td>
<td>$9 \times 10^{18}$</td>
<td>Proton projection imaging, Transmitted Energy Calorimeter</td>
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<tr>
<td>LULI 100 TW</td>
<td>\textsuperscript{65}</td>
<td>50</td>
<td>100-500</td>
<td>$2 \times 10^{20}$</td>
<td>Optical Interferometry</td>
</tr>
<tr>
<td>OMEGA EP</td>
<td>\textsuperscript{64}</td>
<td>750</td>
<td>10-100</td>
<td>$5 \times 10^{18}$</td>
<td>Proton Projection Imaging, Transmitted Energy Calorimeter</td>
</tr>
</tbody>
</table>

Table 1.1: Experiments in channeling performed on different laser systems. $E_L$ is the energy of the channeling laser, $T_L$ is the pulse duration, $n_e$ is the electron density of the plasma.
fraction of energy transmitted through the channel. For the longest channeling pulse durations, up to 80% of the light was transmitted. For the experimental conditions the result is independent of the laser power in the range 100 to 500 GW. Young and Hammer posit that the channeling laser pulse should ramp up adiabatically, that is, in a time longer than it takes for electrons to evacuate the channel. For the parameters of their study, the channel opening time was found to be 100 ps. This time was calculated using a simple equation of motion for the electrons with the ponderomotive force as the driving force.

Another series of experiments on channeling in high density plasma were performed in Japan at the GEKKO-XII (GXII) laser. Takahashi et al. reported that intense laser pulses are capable of evacuating plasma channels beyond the critical surface by a process called superpenetration. X-ray pinhole imaging and x-ray grid image refractometry (GIR) were used to observe the channel penetration beyond critical density plasma. The channeling pulse was 140 J of 1053 nm light delivered in 100 ps in a 30 μm spot with an intensity of $2 \times 10^{17}$ W/cm$^2$. Similar to Young’s experiments, the plasma was generated by ablating a CH foil, which was thicker (100 μm) and was heated by two 351 nm beam with a 100 ps square laser pulse and a laser energy of 10 J. The GXII experiments demonstrate the formation of a large plasma cavity along the channeling beam axis. The x-ray GIR images show that the critical surface of the IR light was pushed closer to the target surface, seemingly coincident with the critical surface of the x-ray radiation, which would indicate a very steep density shelf at the end of the channel. A cone-shaped shockwave was observed that was centered on the front of the channel. These experiments show that intense light is capable of modifying a high density plasma profile. The peak plasma density in the GXII study was higher than that of the Young experiments, because the thick CH target did not completely
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explode before the channeling beam was injected. This led to a plasma with a density gradient normal to the target surface, and closer to what might be expected in FI.

Borghesi et al. performed channeling experiments\(^60\) in preformed inhomogeneous plasma at the Vulcan laser facility. Very thin CH foils (0.1-0.5 \(\mu\)m) were exploded with a frequency-doubled Nd:Glass laser at a wavelength of 527 nm. The plasma forming pulse was 400 ps in duration and had an intensity of \(10^{13}\) W/cm\(^2\). The channeling pulse has a peak intensity of \(9 \times 10^{18}\) W/cm\(^2\) and a duration of 1 ps. The channel was observed using plasma self-emission at the 2nd harmonic (527 nm), a Nomarks interferometer\(^68\) at 527 nm and a Schlieren imaging system. The probe images show a single channel that oscillates in diameter from 5 to 20 \(\mu\)m in diameter\(^60\). The images taken with the Schlieren system show plasma density cavitation in the same region from which the self-emission emanates. The experiments were performed in conjunction with 3-D particle-in-cell simulations, which closely reproduce the observed phenomena, specifically the channel’s transverse and axial shape. The simulations reveal that relativistic electrons traveling with the light pulse serve to stabilize the hosing instability\(^59\) by lowering the refractive index on axis. These experiments show that light is capable of forming a single channel when the pulse is at very high intensity. To increase the intensity of the channeling pulse in these experiments the pulse duration was shorter than in the other channeling experiments (1 ps vs. 10-100 ps).

Experiments on channeling and guiding by the use of a tailored laser focal spot were performed by Durfee and Milchberg\(^69\). Experiments show the guiding of a high-intensity laser pulse through a channel created by the extended focus of an axicon lens. An axicon forms a focus along the propagation direction. Using this technique, high-intensity pulses were guided over many Rayleigh lengths due to the optical guiding nature of the plasma channel. These experiments were conducted in plasmas \(~ 10^{18}\)
cm$^{-3}$ and the scaling to FI-relevant plasmas has not been studied experimentally to date.

OMEGA EP is a new laser system that is capable of conducting FI research. OMEGA EP is ideally suited for this type of experiment, with the capability of generating large, long scale length plasmas with nanosecond UV pulses and simultaneous use of up to two intense short pulse laser beams. Willingale et al. reported evidence of channel formation in low-density plasmas and observed the temporal evolution of laser-instabilities. The channels were observed using side-on proton radiography. The channeling pulse used an energy of 750 J, and a pulse duration of 10 ps that was incident on plasma of density 0.05 $n_c$. The channeling pulse was directed across an ablated plasma plume, which gives a sharp transition from vacuum to plasma, and then a relatively flat density profile for the channeling pulse to interact with. Proton radiography provided images of the evolution of electromagnetic fields with tens of picoseconds time resolution. The total time window of the proton radiography images was 40 ps, which allowed for the observation of the behavior of the channel after the laser pulse. The images show filamentation of the short pulse laser followed by coalescence of the filaments into a single hollow channel. Later, the channel began to form bubble-like structures. The channel had a longitudinal extent of over 1 mm, making it one of the longest channels created in a preformed plasma. The OMEGA EP laser allows scientists to study the channeling physics with the currently highest short-pulse beam energies.

The next logical step on OMEGA EP will be to focus the channeling pulse into an inhomogeneous plasma along a density ramp, which is the topic of this thesis work.
1.6 Thesis Outline

In fast ignition, channeling presents one method of assisting in the ignition energy delivery to a compressed core of an ICF target. The opening time and directional stability are critical to the utility of channeling as an effective means to assist in the delivery of the energy. This thesis will explore channeling in underdense plasma with conditions approaching those expected in a FI target. Chapter 2 discusses the physics of the interaction between light and plasma, and how channels are formed in the plasma. Chapter 3 introduces a new optical diagnostic method for quantitatively assessing the electron density in laser-produced plasma experiments and demonstrates the applicability to channeling experiments. Chapter 4 outlines an experimental program to demonstrate channeling in laser produced plasma using the OMEGA EP laser system at the University of Rochester. Chapter 5 builds on the results of the previous chapter, reporting the results of an experiment to guide a second laser pulse further in the plasma by using a preformed channel. Chapter 6 summarizes the work to date on channeling on OMEGA EP and outlines future modifications to the laser system that may alleviate some of the issues brought to light in the course of this work.
Chapter 2

Physics of the Interaction of High Intensity Light with Matter

Plasma is the most abundant state of matter in the universe. The simplest definition of plasma is an ionized gas. When the electrons are stripped from their atoms, the gas becomes a quasi-neutral system of two mixed fluids. In an ordinary gas, the physical behavior of the gas is governed by collisions between gas atoms or molecules. The statistical average over a large number of individual particles gives the overall macroscopic quantities of density, flow velocity, temperature and heat flux. The difference between gases and plasma is that in plasmas, the particles interact not only by collision but also by the Coulomb force. This allows the particles to screen out fields in a given length scale. For an ionized gas to be considered a classical plasma, three conditions must be met. The field screening length, called the Debye length, must be smaller than the size of plasma. The number of particles within that screened volume must be large so that the statistical average over the distribution is meaningful. The third is that the frequency of electrostatic oscillations due to Coulomb interactions be greater than the collision frequency, so that collective plasma behavior dominates. On Earth however, plasma is the least abundant state of matter owing to the relatively low temperature
and high ionization energy of most matter found in our everyday experience. The Saha equation\textsuperscript{71,72,73,74} calculates the fraction of ionized atoms that can be expected in a plasma in thermal equilibrium, which is given for a hydrogen plasma by

\begin{equation}
\frac{n_i}{n_n} = 6.04 \times 10^{21} \frac{k_B T^{3/2}}{n_i} e^{-E_i/k_B T},
\end{equation}

where $n_i$ and $n_n$ are the number density in units of cm$^{-3}$ of ions and neutral atoms, respectively, $k_B T$ is the temperature of the gas in eV, and $E_i$ is the ionization energy of the gas in eV. The transition from gas to plasma is somewhat arbitrary but typical values of 1% ionization are enough to start showing collective plasma behavior. The fraction of ionized material is noticeable for temperatures $k_B T \sim E_i$. For nitrogen in the atmosphere, this temperature is 160,000 K ($\sim 14$ eV). This high temperature is the reason plasmas are formed only where the energy is high (lightning bolts, nuclear explosions) or where the density is low enough to preclude recombination (aurora, neon lamps). The following sections discuss the basic physics of plasma waves interactions between light and plasma mediated by the dielectric properties of the plasma. The derivations assume that the electrons and ions are fluids and are governed by the conservation laws of continuous media\textsuperscript{75}.

### 2.1 Electrostatic plasma oscillations

The electrons and ions in a plasma form two separate fluids that can move independently of one another and are coupled by electromagnetic forces. Electrostatic fields are generated if significant charge separation occurs. The electrons, because of their smaller mass are pulled back toward the ions to restore local charge neutrality. Their inertia will cause them to overshoot and oscillate back and forth from the electrically
neutral point with a frequency called the electron plasma frequency. This oscillation is so fast that the ions remain relatively motionless due to their higher inertia. In this section, the oscillation frequency of electrons is derived from the electron fluid motion equations. Starting with the momentum balance equation for the electron fluid:

\[
m_{e}n_{e}\left[\frac{\partial v_{e}}{\partial t} + (v_{e} \cdot \nabla) v_{e}\right] = -n_{e}e(E + v_{e} \times B) - \nabla p_{e} - m_{e}n_{e}v_{ei}(v_{e} - v_{i}),
\]

(2.2)

where \(m_{e}\) is the electron mass, \(n_{e}\) is the electron number density, \(v_{e}\) is the electron velocity, \(e\) is the elementary charge, \(E\) and \(B\) are the electric and magnetic fields, \(p_{e}\) is the electron pressure and \(v_{ei}\) is the electron-ion collision rate. Quantities that are written in bold letters are vectors. In the following, the plasma is assumed to be cold \((p_{e} \approx 0)\), non-collisional \((v_{ei} = 0)\), stationary \((v_{e} = 0)\) and absent of external static fields \((E = B = 0)\). All quantities are perturbed about an equilibrium point, and a Taylor expansion up to the first order is used, which results in terms like \(v_{e} = v_{e0} + v_{e1}\) where \(v_{e0}\) is the equilibrium value and \(v_{e1}\) is the perturbed quantity. Then, keeping only terms that are linear in the perturbed value gives

\[
m_{e}n_{e0} \frac{\partial v_{e1}}{\partial t} = -n_{e0}eE_{1}.
\]

(2.3)

This equation can be further simplified by assuming oscillatory solutions of the form

\[
v_{e1} = v_{e1} e^{-i(\omega t - k \cdot x)}.
\]

(2.4)

Applying Equation 2.4 to 2.3 leads to an algebraic expression for \(v_{e1}\) in terms of \(E_{1}\),
\[ -i \omega m_e n_{e0} v_e = -n_{e0}eE_1. \]  
\[ (2.5) \]

Next we can use the continuity equation for the electrons

\[ \frac{\partial n_e}{\partial t} + \nabla \cdot (n_e v_e) = 0, \]
\[ (2.6) \]

and Fourier analyze it to obtain

\[ -i \omega + i k n_{e0} v_{e1} = 0. \]
\[ (2.7) \]

In general, plasmas are assumed to be quasi-neutral but in this specific case for high-frequency waves where the ion inertia is assumed to be infinite there must be charge separation and a resulting electric field. \( E_1 \) and \( n_1 \) are related by Poisson’s equation.

\[ \nabla \cdot E = \frac{e}{\varepsilon_0} (n_{i1} - n_{e1}). \]
\[ (2.8) \]

Subject to the same set of assumptions Eq. 2.8 becomes \( n_e = n_{e0} + n_{e1} \)

\[ i k E_1 = \frac{-e}{\varepsilon_0} n_{e1}. \]
\[ (2.9) \]

Combining the results from Eq. 2.5, 2.7, and 2.9 leads to the dispersion relation

\[ \frac{\omega^2 m_e}{n_{e0}} = \frac{e^2}{\varepsilon_0}, \]
\[ (2.10) \]
which can be rewritten as $\omega^2 = \omega_{pe}^2$ with,

$$\omega_{pe} = \sqrt{\frac{n_e e^2}{m_e \varepsilon_0}} \quad \tag{2.11}$$

the electron plasma frequency, where $\varepsilon_0$ is the permittivity of free space. The frequency $\omega_{pe}$ is a fundamental quantity in plasma physics describing the eigenfrequency of the collective motion of the electron gas against the neutralizing ion background. All perturbed quantities ($v_{e1}, E_1, n_{e1}$) oscillate with the same frequency $\omega_{pe}$. These oscillations do not propagate ($k = 0$) when the plasma has zero temperature. If a finite temperature is assumed, then following the same process leads to the Bohm-Gross\textsuperscript{76,77} dispersion relation,

$$\omega^2 = \omega_{pe}^2 + \gamma_e \frac{k_B T_e}{m_e} k^2, \quad \tag{2.12}$$

where $\gamma_e$ is the adiabatic index of the electron fluid, $T_e$ is the electron temperature and $k_B$ is Boltzmann’s constant. In the warm plasma, the extra thermal contribution to the eigenfrequency of the electrostatic dispersion relation gives a group velocity of

$$\frac{\partial \omega}{\partial k} = \gamma_e \frac{k_B T_e}{m_e \omega_{pe}} \quad \tag{2.13}$$

to the wave and is responsible for the propagation of the electron plasma waves. To find the group velocity, the assumption is made that $\omega_{pe} >> \gamma T_e k / m_e$. 
2.2 Electromagnetic Waves in Plasmas

For a laser-heated plasma, it is important to describe the propagation of electromagnetic waves in plasma. The dielectric properties of the plasma modify the wave. In a plasma the E and B fields in the wave induce sources of current and charge that modify the wave. The response is written in terms of the charge, \( \rho \), and induced current, \( j \), densities:

\[
\rho = \frac{e}{\varepsilon_0} (n_i - n_e),
\]
\[
j = e(n_i \mathbf{v}_i - n_e \mathbf{v}_e),
\]

and Maxwell’s equations can be written as,

\[
\nabla \cdot \mathbf{E} = \rho = \frac{e}{\varepsilon_0} (n_i - n_e),
\]
\[
\nabla \cdot \mathbf{B} = 0,
\]
\[
\nabla \times \mathbf{E} = -\frac{\partial \mathbf{B}}{\partial t},
\]
\[
\nabla \times \mathbf{B} = \mu_0 e(n_i \mathbf{v}_i - n_e \mathbf{v}_e) + \frac{1}{c^2} \frac{\partial \mathbf{E}}{\partial t}.
\]

Taking the curl of Eq. 2.18 yields

\[
\nabla \times (\nabla \times \mathbf{E}) = -\frac{\partial}{\partial t} (\nabla \times \mathbf{B}),
\]

and substituting Eq. 2.19 into Eq. 2.20 then provides the wave equation,
\[ \nabla \times (\nabla \times \mathbf{E}) = -\frac{\partial}{\partial t} \left[ \mu_0 e (n_i \mathbf{v}_i - n_e \mathbf{v}_e) + \frac{1}{c^2} \frac{\partial \mathbf{E}}{\partial t} \right]. \quad (2.21) \]

\[ \left(2.22\right) \]

Invoking the familiar vector triple product identity

\[ \nabla \times (\nabla \times \mathbf{E}) = \nabla (\nabla \cdot \mathbf{E}) - \nabla^2 \mathbf{E}, \quad (2.23) \]

the wave equation then becomes

\[ \frac{1}{c^2} \frac{\partial^2 \mathbf{E}}{\partial t^2} - \nabla^2 \mathbf{E} = -\nabla (\nabla \cdot \mathbf{E}) - \mu_0 e \frac{\partial}{\partial t} (n_i \mathbf{v}_i - n_e \mathbf{v}_e). \quad (2.24) \]

For transverse waves (\( \nabla \cdot \mathbf{E} = 0 \)) this simplifies to

\[ \frac{1}{c^2} \frac{\partial^2 \mathbf{E}^2}{\partial t^2} - \nabla^2 \mathbf{E} = -\mu_0 e \frac{\partial}{\partial t} (n_i \mathbf{v}_i - n_e \mathbf{v}_e). \quad (2.25) \]

Since the light waves are fast compared to the time scale of ion motion, we can assume infinite ion inertia, set \( \mathbf{v}_i = 0 \) and use Eq. 2.3 to solve for \( \mathbf{v}_e \) in terms of \( \mathbf{E}_1 \). The wave equation then becomes

\[ \frac{1}{c^2} \frac{\partial^2 \mathbf{E}}{\partial t^2} - \nabla^2 \mathbf{E} = -\frac{\mu_0 e^2 n_e}{m_e} \mathbf{E} \quad (2.26) \]

The Fourier analysis provides the dispersion relation for an electromagnetic waves in
un-magnetized plasma in its usual form

\[ \omega^2 = c^2 k^2 + \omega_{pe}^2. \] (2.27)

The electron plasma frequency \( \omega_{pe} = \sqrt{n_e e^2/m_e \varepsilon_0} \) and the definition of the speed of light \( c = \sqrt{\mu_0 \varepsilon_0} \) were used to collect the remaining constants. From this expression it is immediately obvious that \( k = 0 \) when \( \omega = \omega_{pe} \). The plasma frequency is a function of density and this creates a characteristic density above which electromagnetic waves of frequency \( \omega_l \) can not propagate in the plasma,

\[ n_c(\omega_l) = \frac{\omega_l^2 \varepsilon_0 m_e}{e^2}, \] (2.28)

where \( n_c \) is referred to as the critical density. Electromagnetic waves do not propagate to densities higher than the critical density. For laser-driven plasma it is often useful to work in terms of the laser wavelength \( \lambda = 2\pi/\omega_l \). Making this substitution, Equation 2.28 can be rewritten as,

\[ n_c = \frac{1.1 \times 10^{21}}{\lambda_{\mu m}^2} \text{cm}^{-3}. \] (2.29)

The wavelength has been normalized to units of \( \mu \text{m} \) for convenience in working with laser wavelengths. For an Nd:Glass laser, the critical density is \( \sim 1 \times 10^{21} \) and for frequency tripled UV drive it increases by a factor of 9. In plasmas that are heated by electromagnetic energy, the critical density marks the location where there is an important change in behavior of electromagnetic waves. The plasma is no longer transparent above this density and the electromagnetic waves do not propagate any further. The electron density is typically specified in fractions of critical density when considering experiments that involve both electromagnetic waves and plasmas.
2.3 Acoustic waves in plasma

In contrast to the high-frequency electrostatic and electromagnetic waves discussed above, plasmas also support longitudinal compression waves as in neutral gases. The difference is that in plasmas, there are few collisions to transmit perturbations as in sound waves due to the screening of potentials in plasma. Sound waves are transmitted from ion to ion through the electric field they generate due to their displacement from the neutral background. The wave is assumed to be neutral because the wave speed is much slower than the electrons response time. The momentum balance equation for the sum of the electron and ion fluid is

\[ m_i n_i \left( \frac{\partial v_i}{\partial t} + \nabla (n_i v_i v_i) \right) = -\nabla P_e - \nabla P_i, \]  

(2.30)

where the electron inertia has been neglected. The electron inertia is considerably smaller than the ion inertia. The continuity equation for the ion fluid is

\[ \frac{\partial n_i}{\partial t} + \nabla (n_i v_i) = 0. \]  

(2.31)

A method similar to that used in the previous sections is employed to linearize these equations about a perturbed equilibrium. The linearized equations become

\[ \frac{\partial n_{i1}}{\partial t} + n_{i0} \nabla v_{i1} = 0 \]  

(2.32)

\[ m_i n_{i0} \frac{\partial v_{i1}}{\partial t} = -\nabla P_{i1} - \nabla P_{e1}. \]  

(2.33)
Taking the gradient of Eq. 2.32 and derivative with respect to time of Eq. 2.33 and solving in terms of $n_{i1}$ yields the wave equation

$$\frac{\partial^2 n_{i1}}{\partial t^2} - \frac{Z\gamma_e k_B T_e + \gamma_i k_B T_i}{m_i} \nabla^2 n_{i1} = 0. \quad (2.34)$$

Here $\nabla P_i$ has been replaced with $\gamma_i T_i \nabla n_i$ by assuming polytropic behavior with $P_i n_i^{-\gamma_i} =$ constant and a similar polytropic equation for the electrons. Here $Z$ is the degree of ionization, and $T_i$ is the ion temperature. Fourier analysis of Eq. 2.34 provides the dispersion relation of ion sound waves in plasma

$$\omega^2 = k^2 \left( \frac{Z\gamma_e k_B T_e + \gamma_i k_B T_i}{m_i} \right) \quad (2.35)$$

and the wave speed is calculated from $\omega/k$ to give

$$c_s = \sqrt{\frac{Z\gamma_e k_B T_e + \gamma_i k_B T_i}{m_i}}. \quad (2.36)$$

The group velocity $\partial \omega / \partial k$ is also $c_s$ so the waves are non-dispersive. The velocity of ion acoustic waves in plasma is also called the ion sound velocity. These waves are akin to sound waves in neutral gas. The speed of the waves is slower than electrostatic or electromagnetic waves in the plasma. The adiabatic index of the electrons is $\gamma_e = 1$ because they transmit energy between each other isothermally on the time scale of ion acoustic fluctuations. Generally the electrons have time to transmit their compression energy to their neighbors during an oscillation period whereas the ions do not. The adiabatic index for the ions is then $\gamma_i = 3$, unless there is time for the energy to be spread isotropically, then $\gamma_i = 5/3$. These waves represent the hydrodynamic motions of the plasma. The nature of the ion acoustic wave is fundamentally different than
sound waves in gases. In an ion acoustic wave, it is the temperature of the electron gas that provides the restoring force and the mass of the ions that carry the inertia of the wave. Typically the electrons are hotter than the ions \(T_e \gg T_i\). In some cases the electrons are cold \(T_e = 0\), and the expression is identical to ordinary pressure waves in gases with

\[
c_s = \sqrt{\frac{\gamma k_B T_i}{m_i}}. \tag{2.37}
\]

### 2.4 Generation of Long Scale Length Plasmas

A laser beam propagating in plasma can be absorbed by a process called collisional absorption or inverse bremsstrahlung absorption. Free electrons oscillate in the laser field and the free electron gas is heated through Coulomb collisions between electrons and ions. The collisions transfer energy from the laser beam to the electrons. The collisions can heat and ionize neutral atoms, releasing more electrons in a process called avalanche ionization. As a consequence the material is heated to a high temperature in the laser spot.

Starting from the electron momentum equation, this time keeping the momentum transfer term \(m_e n_e v_{ei}(v_i - v_e)\),

\[
m_e n_e \left[ \frac{\partial v_e}{\partial t} + (v_e \cdot \nabla) v_e \right] = -n_e e(E + v_e \times B) - \nabla p_e - mn_e v_{ei}(v_e - v_i), \tag{2.38}
\]

and considering the laser irradiation as a high frequency electric field

\[
E(t) = E_0 e^{-i\omega_l t} \tag{2.39}
\]

where \(\omega_l\) is the angular frequency of the laser. In Sec. 2.1 we neglected the collision
frequency $v_{ei}$. Now we keep the effect of collisions to obtain an expression for the
collisional heating of a plasma. We assume $v_0 = E_0 = B_0 = 0$ and keep the first order
terms to obtain,

$$\frac{\partial v_{e1}}{\partial t} = -\frac{e}{m_e} E_1 - v_{ei} v_{e1}. \quad (2.40)$$

Since the forcing field from the laser is time-harmonic we will use the plane wave
assumption to arrive at

$$v_{e1} = -\frac{-ie E_1}{m_e (\omega_l + iv_{ei})}. \quad (2.41)$$

This expression can be used in the dispersion relation found in the earlier section to
yield a new dispersion relation including collisions,

$$\omega^2 = c^2 k^2 + \frac{\omega_{pe}^2}{1 + iv_{ei}/\omega}, \quad (2.42)$$

which can be simplified by invoking the binomial approximation assuming $v_{ei}/\omega << 1$,

$$\omega^2 = c^2 k^2 + \omega_{pe}^2 (1 - \frac{iv_{ei}}{\omega}). \quad (2.43)$$

The corresponding dielectric function, $\varepsilon$, including collisions is found to be

$$\varepsilon = 1 - \frac{\omega_{pe}^2}{\omega^2 + v_{ei}^2} \left(1 + \frac{v_{ei}}{\omega}\right). \quad (2.44)$$

The complex refractive index, $N$, is obtained by taking the square root of $\varepsilon$. 
\[ N = n - i\tilde{n} = \sqrt{\varepsilon} \] (2.45)

with

\[ n = \frac{1}{\sqrt{2}} \sqrt{A + \sqrt{A^2 + B^2}}, \] (2.46)
\[ \tilde{n} = \frac{1}{\sqrt{2}} \sqrt{-A + \sqrt{A^2 + B^2}}, \] (2.47)

where

\[ A = 1 - \frac{\omega_{pe}^2}{v_{ei}^2 + \omega^2}, \] (2.48)
\[ B = \frac{v_{ei} \omega_{pe}^2}{\omega (v_{ei}^2 + \omega^2)}. \] (2.49)

The complex propagation coefficient, also called the wavenumber, is

\[ \gamma = \alpha + i\beta \] (2.50)

with

\[ \alpha = \tilde{n} \frac{\omega}{c}, \] (2.51)
\[ \beta = n \frac{\omega}{c}. \] (2.52)

When the collision rate is low \(v_{ei} << \omega\), and \(\omega_{pe} < 1/2\omega\) the real part of the refractive index is the same as without collisions.
\[ n = \sqrt{1 - \frac{\omega_{pe}^2}{\omega_l^2}} \]  

while the imaginary part of the refractive index is given by \( \tilde{n} = \frac{1}{2}(\nu_{ei}/\omega_l)(\omega_{pe}/\omega_l)^2/\sqrt{(1 - \omega_{pe}^2/\omega_l^2)} \). The absorption coefficient \( \kappa_{IB} \) of the laser radiation in the plasma medium is related to the imaginary part of the refractive index \( \tilde{n} \) according to \( \kappa_{IB} = 2(\omega/c)\tilde{n} \), which results in

\[ \kappa_{IB} = \frac{\nu_{ei}}{c}\left(\frac{n_e}{n_c}\right) \frac{1}{\sqrt{(1 - \frac{n_e}{n_c})}} \]  

(2.54)

using \( (\omega_{pe}/\omega_l)^2 = n_e/n_c \). Since the collision frequency \( \nu_{ei} \propto Zn_e/T_e^{3/2} \), the absorption coefficient depends on electron density, ionization degree, and temperature in the following way

\[ \kappa_{IB} \propto \frac{Zn_e^2}{T_e^{3/2}} \frac{1}{\sqrt{(1 - \frac{n_e}{n_c})}} \]  

(2.55)

Collisional absorption is the dominant absorption mechanism for laser intensities up to \( \sim 10^{15} \) W/cm². The process describes the absorption of electromagnetic radiation in the presence of electron-ion collisions. This explains the term inverse bremsstrahlung, because bremsstrahlung is the emitted electromagnetic radiation as the result of electron-ion (and electron-electron) collisions.

The absorbed laser energy heats the plasma and causes it to expand. The coronal plasma has a high thermal conductivity and is therefore close to isothermal. An isothermal gas that expands in a self-similar manner is governed by the self-similarity variable
\( r/t \), where \( r \) is the 1-D expansion distance, and \( t \) is the time the plasma has had to expand. The equation for a freely expanding plasma can be found by starting from the continuity and force equations for the ion fluid given by

\[
\frac{\partial n}{\partial t} + \frac{\partial}{\partial r}(nv) = 0 \quad (2.56)
\]

\[
\frac{\partial v}{\partial t} + v \frac{\partial v}{\partial r} = -c_s^2 \frac{1}{n} \frac{\partial n}{\partial r} \quad (2.57)
\]

The self-similar solution is found by choosing \( n = f(\xi) \) and \( v = g(\xi) \) where \( f \) and \( g \) are functions to be determined and \( \xi \) is the self-similarity variable \( r/t \). Substituting into Eqs. 2.56 and 2.57 gives the coupled non-linear differential equation pair

\[
\frac{d}{d\xi} f(g - \xi) + f \frac{dg}{d\xi} = 0, \quad (2.58)
\]

\[
\frac{d}{d\xi} g(g - \xi) + c_s^2 \frac{df}{d\xi} f = 0. \quad (2.59)
\]

These two equations can be separated to yield

\[
g = \xi + c_s \quad (2.60)
\]

\[
\frac{1}{f} \frac{df}{d\xi} = -\frac{1}{c_s} \quad (2.61)
\]

and the substituting back \( \xi = r/t \), \( v = g(\xi) \) and \( n = f(\xi) \) yields

\[
v = c_s + \frac{r}{t} \quad (2.62)
\]

\[
n = n_0 e^{-\frac{r}{c_st}} \quad (2.63)
\]

The expansion speed is the ion-acoustic velocity, \( c_s \). The plasma density is profile is a function of the distance \( r \) and the time, \( t \), after the start of the expansion. In the
freely expanding plasma there is a characteristic length scale $L_s = \frac{c_s t}{\kappa}$ for the density inhomogeneity. The local density scale length is defined as $L_s = (\nabla n_e / n_e)^{-1}$. As time increases, the density becomes increasingly homogeneous. Two-dimensional hydrodynamic simulations using the code DRACO show that the maximum achievable scale length in planar geometry is obtained by maximizing the overlapped-laser-beam intensity while providing enough time for the plasma to reach the steady state. For the experimental conditions encountered on OMEGA EP, the scale length reaches a steady state after about 1.5 ns, and the asymptotic scale length is given by $L_s = 250 \mu m (I_{14})^{1/4}$.79

2.5 Ponderomotive Force

Under circumstances in our normal daily life, radiation exerts a small pressure on surfaces it is incident upon. The radiation pressure exerted on surfaces is typically too small to be measured. When the intensity of the light is much higher, as is the case with high powered laser systems, the light pressure can be equal to or greater than the surrounding ambient pressure. When the light pressure is considered in the context of plasma physics, the resulting force arising from gradients in the light pressure is termed ponderomotive force. The ponderomotive force is responsible for electromagnetic waves coupling to plasmas in all parametric instabilities. In the next section, the ponderomotive force is derived from the motion of charged particles in an electromagnetic wave.

2.6 Physics of Channeling

The process of channeling is due to a force exerted on electrons by the laser called ponderomotive force. The ponderomotive force is felt by electrons (and to a much
lesser extent by the ions) when oscillating in a spatially inhomogeneous electromagnetic wave. The force can be quite large when the plasma is exposed to the high intensity laser fields created by chirped pulse amplification lasers. For instance, on a modern high-power laser system like OMEGA EP, the ponderomotive pressure can reach Gbar levels, and can overwhelm the hydrodynamic pressure in the coronal plasma. In this section, we derive the physics of ponderomotive force for an arbitrary EM wave and discuss the implications of high intensity laser pulses in plasmas.

2.6.1 Relativistic Ponderomotive Force

The most general treatment of ponderomotive potential and ponderomotive force on a free charge can be derived using a cycle-averaging technique in phase-space. The benefit of this method compared to the standard perturbative analysis is that it dispenses with a priori assumptions of time-harmonic fields and small oscillation amplitude. The latter assumption is not appropriate when considering modern high power laser systems. The relativistic Lagrangian of a single electron charge, $e$ in an electromagnetic field specified by a vector potential $A$ can be constructed as

$$L(x, v, t) = -\frac{mc^2}{\gamma} + ev\cdot A - e\Phi,$$  \hspace{1cm} (2.64)$$

where $x$ is the space coordinate, $v$ is the velocity, $t$ is time, $\gamma = \sqrt{1 - \frac{v^2}{c^2}}$, $m$ is the electron rest mass, $e$ is the elementary charge, and $\Phi$ is the scalar potential. We consider here the transverse gauge where $\nabla \cdot A = 0$. To separate the time scale of fast oscillatory motion of the charge and the slow drift of the oscillation center, we introduce an action-angle variable pair $S(x, t)$ and $\psi(x, t)$. $S$ is the action and is defined as
\[ S = \int_{\psi_1}^{\psi_2} L(x, v, t) \frac{dt}{d\psi} d\psi. \] (2.65)

Under a change from one inertial coordinate system to another, \( S \) and \( \psi \) are invariant. Because of the Lorentz invariance of \( S \) and \( \psi \), the Lagrangian, given by

\[ \mathcal{L}_0 = L(\psi) \frac{dt}{d\psi}, \] (2.66)

is also invariant under change of inertial system. We construct a cycle-averaged Lagrangian \( \mathcal{L}_0 \) assuming that the motion of the charge is periodic (not necessarily harmonic) and \( \psi \) is normalized to \( 2\pi \),

\[ \mathcal{L}_0 = \frac{1}{2\pi} \int_{\psi}^{\psi + 2\pi} L(x_0', v_0', t(\psi')) \frac{dt}{d\psi} d\psi', \] (2.67)

where \( x_0, v_0 \) are the coordinates of the center of oscillation of the charge. The equations of motion for the oscillation center can be found by the Lagrange equations of motion

\[ \frac{d}{dt} \frac{\partial L_0}{\partial v_0} - \frac{\partial L_0}{\partial x_0} = 0, \] (2.68)

where \( L_0 = \mathcal{L}_0 d\psi/dt \). It can be shown that Equation 2.67 representing the motion of the oscillation center holds for adiabatic (slow compared to oscillation time) changes in the system and consequently the Hamiltonian representing the total cycle-averaged energy of the charge is an adiabatic invariant.\(^8\)

If we consider a monochromatic wave of arbitrary strength in vacuum, we may
choose $\Phi = 0$. When we introduce the effective mass as $m_{\text{eff}} = \mathcal{L}_0 \gamma_0 (d\psi/dt)/c^2$, the cycle-averaged Lagrangian and Hamiltonian are connected to a single free particle with an effective mass that varies in space and time,

$$L_0(x_0, v_0, t) = \frac{-m_{\text{eff}}c^2}{\gamma_0},$$

$$H_0(x_0, p_0, t) = \gamma_0 m_{\text{eff}}c^2,$$  

$$\gamma_0 = \left(1 - \frac{v_0^2}{c^2}\right)^{-1/2},$$  

$$p_0 = \gamma_0 m_{\text{eff}}v_0.$$  

These expressions are suitable for any electromagnetic field where the charge has a distinct oscillation center. For an electromagnetic wave the effective mass is given by

$$m_{\text{eff}} = m_e \left(1 + \frac{e^2 |A|^2}{\alpha m_e^2 c^2}\right)^{1/2}$$

where $\alpha = 2$ for linear polarization and $\alpha = 1$ for circular polarization. We choose $\alpha = 2$ and an inertial frame where $v_0 = 0$ and the ponderomotive force is found using Eq. 2.68 to be,

$$F_p = \frac{\partial L_0}{\partial x_0} = -c^2 \nabla m_{\text{eff}} = -c^2 m_e \nabla \gamma = -m_e c^2 \nabla \left(1 + \frac{e^2 |A|^2}{2m_e^2 c^2}\right)^{1/2}.$$  

This result yields several insights into the ponderomotive force acting on charges. The ponderomotive force of an electromagnetic wave of any strength can be found through cycle-averaging the Lagrangian of the center of motion. Cycle-averaging must be done over the invariant phase, because averaging over $t$ becomes incorrect as the oscillation velocity approaches relativistic speeds. The oscillation center must be defined by Lorentz-invariant variables (in our case, $\psi(x, t)$). We can express the ponderomotive
force more compactly by normalizing the laser vector potential by
\[ a = eA/(m_ec) \]
and
the expression for ponderomotive force of linear polarized light is then

\[ F_p = -m_ec^2\nabla\left(1 + \frac{|a|^2}{2}\right)^{1/2}. \quad (2.75) \]

### 2.6.2 Steady State Channeling for a Fixed Laser Intensity

In the last section we showed that the ponderomotive force is proportional to the gradient of the laser intensity. The ponderomotive force of a laser beam with a Gaussian transverse profile ejects electrons out of the focal spot. This sets up a large space charge force which ejects the ions radially from the laser axis, called a Coulomb explosion. A simple physics model of the particle dynamics in the focal spot of an intense laser can be constructed using the expression for ponderomotive force, and the ion-acoustic wave equation in plasma. The electron motion is described by a radial force balance,

\[ e\nabla_r \phi = m_ec^2\nabla_r\left(1 + \frac{a^2}{2}\right)^{1/2} + \frac{1}{n_e}\nabla_r P_e, \quad (2.76) \]

where \( e\nabla_r \phi \) is the space charge force, \( m_ec^2\nabla_r\left(1 + \frac{a^2}{2}\right)^{1/2} \) is the ponderomotive force, \( P_e = n_eT_e \) is the electron pressure, and \( n_e \) is the electron density. The ion motion is described by the continuity equation

\[ \frac{\partial \delta n_i}{\partial t} = -n_i\nabla_r \cdot \mathbf{v}_r, \quad (2.77) \]

and the momentum equation

\[ \frac{\partial \delta \mathbf{v}_ri}{\partial t} = -Ze\nabla_r \phi, \quad (2.78) \]
where $\delta n_i$ and $n_i$ are the perturbed and ambient ion densities, $v_{ri}$ is the ion velocity in the radial direction, and $T_i << T_e$ is assumed. Combining these equations yields the channeling density cavitation that can be calculated starting from the ion-acoustic equation driven by ponderomotive force given by

$$
\left( \frac{\partial^2}{\partial t^2} - c_s^2 \nabla^2 \right) \frac{\delta n_i}{n_i} = \frac{Z m_e c^2}{m_i} \nabla^2 \left( 1 + \frac{|a|^2}{2} \right)^{\frac{1}{2}},
$$

(2.79)

where $\delta n_i$ is the density excursion from equilibrium ($\delta n_i = n_i/n_i$), $Z$ is the ionization state of the plasma, $m_e$ and $m_i$ are the electron and ion masses respectively and $a$ is the normalized laser vector potential. If the laser pulse envelope is slowly varying with respect to the laser oscillatory period then we drop the component of the Laplacian in the direction of the laser pulse $\nabla^2_z$ and focus only on the radial component $\nabla^2_r$. For the sake of analytic simplicity, we assume the non-relativistic case with $a^2 < 1$. The equation becomes

$$
\left( \frac{\partial^2}{\partial t^2} - c_s^2 \nabla^2_r \right) \frac{\delta n_i}{n_i} = \frac{Z m_e c^2}{4 m_i} \nabla^2_r a(r,t) \left( r, t \right)^2,
$$

(2.80)

where $c_s = \frac{c}{\sqrt{Z}}$. Equation 2.80 has a solution\(^8^3\) of the form

$$
\frac{\delta n_i}{n_i} = \frac{1}{2 c_s} \int_{0}^{t} \int_{r - c_s(t-t')} \left[ \int_{r + c_s(t-t')} \frac{Z m_e c^2}{4 m_i} \nabla^2_r a(r', t') dr' \right] dt'.
$$

(2.81)

Using Eq. 2.81, we calculate the residual density in an evacuated channel by using the OMEGA EP laser system. The channeling beam on OMEGA EP can be up to 2 kJ with a pulse width of 10 - 100 ps and a best focus spot of 20 $\mu$m radius, which yields an on target intensity of $5 \times 10^{18}$ W cm$^{-2}$ for the 100 ps case. For the 100 ps pulse, the normalized vector potential, $a_0$ is 0.85. If we consider CH plasma with $T_e = 1$ keV and
Figure 2.1: This image shows the on-axis cavitation due to an intense laser pulse at two times: 30 ps into the pulse (top) and 80 ps into the pulse (bottom). The red dashed curve indicates the background plasma density without the channel. This channel was dug using a 100 ps, 2 kJ laser pulse in a 1 keV CH plasma, similar to the conditions found on OMEGA EP planar foil explosions.
Z = 3.5, then the transverse density profile predicted by Eq. 2.81 is shown in Figure 2.1 at two times in the pulse (30 and 80 ps).

2.7 Light Pressure from Ponderomotive Force

If we consider the steady-state behavior of an electromagnetic plane wave reflecting from an overdense interface, we can recover the expression for the light pressure starting from the expression derived for the ponderomotive force. In this example we take a standing plane wave in a plasma of the form

\[ E(x, t) = E_0 + E_r = E_0 e^{ikx - i\omega t} + E_r e^{-ikx - i\omega t}, \]  
\[ E^*(x, t) = E^*_0 + E^*_r = E^*_0 e^{-ikx + i\omega t} + E^*_r e^{ikx + i\omega t}, \]

where \( E_0 \) and \( E_r \) are the incident and reflected wave respectively and \( E^* \) denotes the complex conjugate of \( E \). The electromagnetic wave must obey the dispersion relation of the plasma (ignoring collisions, see Eq. 2.27) so we can set up the stationary wave equation for \( E \) and \( E^* \)

\[ \nabla^2 E + k_0^2 (1 - \frac{\omega_p^2}{\omega^2}) E = 0, \]  
\[ \nabla^2 E^* + k_0^2 (1 - \frac{\omega_p^2}{\omega^2}) E^* = 0, \]

where \( k_0 \) is the wavenumber in vacuum \( k_0 = \frac{\omega}{c} \). Considering normal incidence along the x-direction the pressure of the light exerted on a surface can be written as the sum of the ponderomotive force in the direction of propagation,

\[ p_L = \int_{x>0} f_p dx = \int_{x>0} \frac{E_0^2 \omega_p^2}{4 \omega_r^2} \frac{\partial}{\partial x} (EE^*) dx, \]
where we use Eq 2.11 to express \( f_p \) in terms of \( E, \omega_l, \omega_{pe}, \) and \( \varepsilon_0 \). Here \( \varepsilon_0 \) is the free space permittivity, \( \omega_{pe} \) is the local plasma frequency and \( \omega_l \) is the laser frequency. In section 2.6.1, the ponderomotive force was derived on a per particle basis, so the force must be multiplied by the local plasma density \( n_e \) to get the bulk force. We can expand the derivative in Eq. 2.86 to

\[
\frac{\partial}{\partial x}(EE^*) = \frac{\partial E}{\partial x}E^* + \frac{\partial E^*}{\partial x}E = \frac{\partial}{\partial x}(EE^* + \frac{1}{k_0} \frac{\partial E}{\partial x} \frac{\partial E^*}{\partial x}) \tag{2.87}
\]

by using Eqs. 2.84 and 2.85. First, the Eq. 2.84 and 2.85 are multiplied by \( \frac{\partial E^*}{\partial x} \) and \( \frac{\partial E}{\partial x} \) respectively. This yields

\[
\frac{\omega_p^2}{\omega_l^2} \frac{\partial E^*}{\partial x} = \frac{1}{k_0^2} \frac{\partial^2 E}{\partial x^2} \frac{\partial E^*}{\partial x} + E \frac{\partial E^*}{\partial x}, \tag{2.88}
\]

and

\[
\frac{\omega_p^2}{\omega_l^2} \frac{\partial E}{\partial x} = \frac{1}{k_0^2} \frac{\partial^2 E^*}{\partial x^2} \frac{\partial E}{\partial x} + E^* \frac{\partial E}{\partial x}. \tag{2.89}
\]

Recognizing the left hand side of these equations are the integrand in equation 2.86 it may be rewritten as

\[
p_L = \frac{\varepsilon_0}{4} \int_{x>0} f_p dx \tag{2.90}
\]

\[
= \frac{\varepsilon_0}{4} \int_{x>0} \left[ \frac{1}{k_0^2} \frac{\partial^2 E}{\partial x^2} \frac{\partial E^*}{\partial x} + E^* \frac{\partial E}{\partial x} + \frac{1}{k_0^2} \frac{\partial^2 E^*}{\partial x^2} \frac{\partial E}{\partial x} + E \frac{\partial E^*}{\partial x} \right] dx. \tag{2.91}
\]

This equation can then be expressed as a derivative

\[
p_L = \frac{\varepsilon_0}{4} \int_{x>0} f_p dx = \frac{\varepsilon_0}{4} \int_{x>0} \frac{\partial}{\partial x} [EE^* + \frac{1}{k_0^2} \frac{\partial E}{\partial x} \frac{\partial E^*}{\partial x}], \tag{2.92}
\]
and evaluated to find

\[ p_L = \frac{\varepsilon_0}{4} \left[ (E_0 + E_r)(E_0^* + E_r^*) + (E_0 - E_r)(E_0^* - E_r^*) \right] \]  

(2.93)

\[ = \frac{\varepsilon_0}{2} (|E_0|^2 + |E_r|^2). \]  

(2.94)

The expression for light pressure is recovered by recognizing that the incoming light intensity is \( I = \varepsilon_0 E_0^2 / 2 \) and the reflectivity is the ratio \( R = E_r^2 / E_0^2 \). The expression above becomes

\[ p_L = (1 + R) \frac{I}{c}. \]  

(2.95)

From here we see that light pressure emerges as a consequence of ponderomotive force in the direction of the laser beam. Gradients in plasma density in the direction of the plane wave will cause forces to develop in the forward direction which are measured as pressure on the surface from which the radiation is reflected.

### 2.8 Hole-boring

The light impinging on an absorbing or reflecting interface imposes pressure as was described in the previous section and might push the vacuum-plasma boundary forward with sufficiently high pressure. A simple model is considered that describes the light reflection from a moving interface and the momentum transfer. In a frame moving with the material interface, the light pulse is considered to be a “macro-particle” colliding with the plasma. Material does not leave the boundary, but is reflected from the interface. The momentum is conserved over the collision, given by

\[ P_L + P_L' = P_p + P_p', \]  

(2.96)
2.8. HOLE-BORING

where \( P_L, P'_L \) and \( P_p, P'_p \) are the momenta of the light and plasma before and after the collision respectively. This leads to that the forward directed light pressure is balanced with the dynamic pressure generated by the head of the channel pushing plasma out of the way. The model was put forth by Kruer et al. initially to describe density steepening in underdense plasmas\(^{14}\). With the invention of high-powered lasers capable of radically modifying the density surface with light pressure, the model was extended by Wilks et al. to describe the recession velocity of a solid target surface irradiated by an intense laser pulse\(^{5}\). Figure 2.2 shows schematically the balance of pressures acting on the boundary. The left side indicates the pressure from the laser light. The right side represents the pressure arising from plasma entering the reflection front. The balance of pressures can be written as

\[
\frac{I_L}{c} (1 + R) = \frac{2n_e m_i}{Z} v_c^2
\]

where \( v_c \) is the interface’s forward going velocity, \( n_e \) is the local electron density ahead of the interface, \( m_i \) is the ion mass and \( Z \) the charge state. The velocity increases for increasing laser intensity, and slows for an increasing density. The velocity can then be found as,

\[
v_c = \sqrt{\frac{I_L Z (1 + R)}{2n_e m_i c}}. \tag{2.98}
\]

This is called the “hole-boring” velocity because the model describes how an intense laser pulse pushes into a plasma interface with a steep density gradient.
Figure 2.2: A simple model of hole-boring using light in a plasma. The light pressure on the left side is balanced against the dynamic pressure arising from the outward flow on the right side. Figure adapted from\textsuperscript{14}.

2.9 Summary

This chapter introduced the definition of plasma and the physical basis of the oscillations and waves in plasma were derived. The interaction between light and plasma was discussed, with special emphasis on the secular forces arising from the interaction. Finally, a specific example of the forces interacting with plasma was presented that ultimately led to the modification of the plasma density profile. Light pressure as a secular force is shown to emerge from the formal definition of ponderomotive force. Channeling, the radial evacuation of plasma from the focal spot of the laser beam, and hole boring, the pushing of an overdense surface by light pressure are shown to arise from the ponderomotive force of the laser beam. The physical demonstration of these concepts is shown in the experimental section of this thesis.
Chapter 3

Experimental Technique: Angular Filter Refractometry

A novel diagnostic technique based on refraction was developed to observe the modification of the plasma by the channeling beam. Angular Filter Refractometry (AFR)\textsuperscript{86} can probe high-density, long-scale-length plasmas relevant to high-energy-density physics. AFR measures electron plasma densities up to 10\textsuperscript{21} cm\textsuperscript{-3} by using a 263-nm probe laser\textsuperscript{15,86}. This chapter describes the development and calibration of AFR. This includes the calibration of the angular filters against a known reference and measurements of the optical resolution of the system. A Fourier optical model of AFR is presented to reproduce the diffraction features seen in the AFR image. The chapter concludes with example results of an experiment to show the utility of this diagnostic.

3.1 Fourth-harmonic Optical Probe

A 10-ps, 263-nm (4\omega) laser was built to probe plasmas produced on the OMEGA EP Laser System. Figure 3.1 shows a schematic of the optical probe system from the laser source to the detectors. The laser system starts with a 10-ps, 1053 nm oscillator.
The system uses a large-aperture ring amplifier (LARA)\(^8\) to amplify the pulses to an energy of 100 mJ per pulse. The light is optically relayed to a table located at the target chamber wall, where the light is frequency quadrupled using two KD*P crystals\(^8,8^9\) producing pulses with a wavelength of 263 nm and an energy of 10 mJ. The light is focused by an f/25 lens into the vacuum chamber 8.5 cm in front of the center of the chamber. The light passes through the plasma located in the center of the chamber and is collected by an f/4 telescope and relayed over a distance of \(\sim 5 \) m to the diagnostics for imaging. The collected signal is split into various imaging channels. The 263-nm wavelength was chosen to allow optical probing of high densities and to discriminate against scattered light from the high energy 351-nm wavelength drive laser. The 10-ps pulse width minimizes temporal blurring due to hydrodynamic motion. A narrow (\(\pm 2 \) nm) bandpass-filter around 263 nm passes the probe light and efficiently rejects the background bremsstrahlung emission from the plasma.

Figure 3.2 shows a mechanical rendering of the catadioptric telescope designed to image the object from the center of the vacuum chamber to a location outside the chamber wall. The light is collimated by a five-element achromat. The collection telescope operates at f/4, giving a diffraction limited resolution of \(\sim 1 \mu m\). The majority
of the optical power in the system is provided by the meniscus lens. The meniscus cavity, the space between a tilted flat plate and meniscus lens (Figure 3.2), acts as a powered reflector for the first pass through it. After reflection off the first surface of the meniscus, the light reflects off a plane surface tipped by 4.7 degrees. The beam then passes through the meniscus lens on the way to the output port. The curvature of the second surface of the meniscus corrects the spherical aberration in the optical system. A single meniscus lens can be used to correct spherical aberration. The collector is similar to a Maksutov telescope where the first meniscus surface acts as the primary mirror and the second surface acts as a corrector. This arrangement allows for a near diffraction limited, achromatic image to be produced outside the target chamber. The spherical aberrations arising from the propagation through a protective blast window and the EP chamber vacuum windows are compensated by design in the curvature of the rear surface of the meniscus. The meniscus is flexured about the center of the first surface to allow pivoting to correct for shear between the two surfaces of the meniscus lens during fabrication as the optical design is sensitive to shear and has no way of correcting for it.

Figure 3.3 shows a mechanical rendering of the whole optical system including the target chamber. The probe laser propagates inside the gold colored reentrant vacuum
tube shown in the lower left and, passes through the center of the target chamber and is collected by the silver colored tube in the upper right.

The image provided outside the target chamber is relayed to an optical bench located a distance away from the target chamber. Figure 3.4 shows a schematic of the optical layout on the diagnostic bench. Upon emerging from the relay periscope, the light is split into 4 channels. Two channels are dedicated to the pointing and centering of the beam, denoted here by the red rays. Images of the near field and the far field of the beam are measured with GigE manta CCD cameras\textsuperscript{92}, using a magnification of \(4.1 \times\). The two scientific imaging legs are shown by the blue rays, which are referred to as AFR1 and AFR2 and capture a magnified image of the object plane at TCC with scientific PIXIS cameras\textsuperscript{93}. The cameras contain 2048 x 2048 pixels with a pixel size of 13.5 \(\mu\)m. The imaging legs AFR1 and AFR2 have a magnification of 7.3 \(\times\) and 7.1 \(\times\), respectively.

The resolution of the imaging system was measured by taking the 90-10 \% intensity width of a static grid target. A fine copper grid with 60 \(\mu\)m grid spacing was placed in the beam path at target chamber center and an image was recorded on the AFR1 and AFR2 cameras. Figure 3.5(a) shows an image of the grid and in Fig.3.5(b), a horizontal lineout averaged over the central grid region. Figure 3.5(c) shows the transmission edge function. The resolution of the imaging system is found to be 3.6 \(\pm\) 1.8 \(\mu\)m in the object plane over the measurable field-of-view.

The probe rays are collimated and pass through a bandpass filter with an extinction ratio of \(10^5\) for wavelengths outside of 263\(\pm\) 2 nm. This is necessary because the heated plasma emits a broad spectrum of light in the UV region. Since the signal beam is weak (10 mJ) compared to the drive lasers (> 1 kJ), significant bandpass filtering at the signal frequency is required.
3.1. FOURTH-HARMONIC OPTICAL PROBE

Figure 3.3: Mechanical layout of the $4\omega$ probe laser system. The laser beam is generated on an optical table at the bottom of the image, propagated through tubes then frequency converted to the 4th harmonic and then is focused into the target chamber. The f/4 optical collection telescope in the upper right collects the refracted light from the experiment. The light is down collimated, filtered through a bandpass filter, and relayed to an optical bench for optical analysis (upper right).

Figure 3.4: Layout of the $4\omega$ optical diagnostic bench. The image from the object located at target chamber center is relayed to the optical bench shown here. The beam is split into independent legs which allow for simultaneous optical measurements.
Figure 3.5: (a) Image on the CCD detector of a wire grid positioned in the center of the $4\omega$ beam in target chamber center. (b) Horizontal lineout through the grid image. (c) The 90-10 % width (3.6 ± 1.8 μm) gives a measure of the minimum resolvable features in the object plane.
The next section discusses the physical basis of probing plasma objects with an optical probe.

### 3.2 Optical Probing of Plasma

In section 2.4, the local refractive index, \( n \), of a plasma was shown to be directly related to the local electron density \( n_e \). By measuring the accumulated phase along the optical path of a probing beam passing through the transparent, refractive plasma, a measurement of the electron density is possible. The techniques of optical probing of plasma rely on the differential deviation of parts of the probe beam due to refraction. The chord-integrated phase of the probe beam can be related to the density using Equation 2.46 for the refractive index of a plasma,

\[
\Phi(x, y) = \int_{-\infty}^{\infty} k_p \, dl
\]

\[
= \frac{\omega_l}{c} \int_{-\infty}^{\infty} n \, dl
\]

\[
= \frac{\omega_l}{c} \int_{-\infty}^{\infty} \sqrt{1 - \frac{n_e(x, y, z)}{n_c}} \, dl
\]

\[
\Phi(x, y) \approx \frac{\pi}{\lambda_p n_{cr}} \int_{-\infty}^{\infty} n_e(x, y, z) \, dl, \text{ for } n_e << n_c
\]

where \( \lambda_p \) is the probe laser wavelength, \( \Phi(x, y) \) is the total accumulated phase of a probing ray as a function of the image plane coordinates, \( k_p \) is the wave number of the probe beam, \( \int dl \) is the total path that the ray propagates through the plasma, \( n_e(x, y, z) \) is the plasma density, and \( n_{cr} \) is the critical plasma density for the probe laser given by Eq. 2.28, which is \( 1.6 \times 10^{22} \, \text{cm}^{-3} \) for \( \lambda_p = 263 \, \text{nm} \). Equation 3.4 shows that the total accumulated phase is proportional to the chord integrated electron density along the
CHAPTER 3. EXPERIMENTAL TECHNIQUE: ANGULAR FILTER REFRACTOMETRY

Figure 3.6: In this schematic the probing geometry of a ray passing through a plasma is shown. The ray enters at an impact parameter of \( y \).

path through the plasma. By making the assumptions that the plasma is axis-symmetric about the \( y \)-axis and the optical path is straight along the \( z \)-direction (see Figure 3.6), this equation can be Abel inverted\(^{94}\) along the \( x \)-axis for a given \( y \) to solve for \( n_e \) as a function of the accumulated phase. This assumption has been demonstrated to be accurate to \( 2 \times 10^{19} \) cm\(^{-3} \) in plasma density so long as the following conditions are satisfied: (i) The density probed by the probing radiation should be ten times less than the critical density for the probe, (ii) Strongly refracted rays (\( > 10^{9} \)) are rejected by the aperture stop of the optical system, and (iii) the density is decreasing as a function of impact parameter\(^{95}\).

The schematic in Fig. 3.6 shows the geometry. The plasma extends to a certain radius \( R \) where the plasma density approaches zero. The phase can then be written as

\[
\Phi = \frac{\pi}{\lambda_p n_{cr}} \int_{-\sqrt{R^2-y^2}}^{\sqrt{R^2-y^2}} n_e[r(z)] dz, \tag{3.5}
\]
which by changing variables may be rewritten as an Abel transform

\[
\Phi = \frac{2\pi}{\lambda_p n_{cr}} \int_y^R \frac{n_e(r)r}{\sqrt{r^2 - y^2}} dr.
\]  

(3.6)

The Abel inversion formula is well known and is given by\textsuperscript{94} (Appendix A),

\[
n_e(r) = -\frac{\lambda_p n_{cr}}{\pi} \int_r^R \frac{d\Phi}{dy} \frac{1}{\sqrt{y^2 - r^2}} dy.
\]  

(3.7)

The local refraction angle refraction \((\theta(x, y))\) of a probing ray is related to the transverse gradient of the phase of the wavefront at the point \(x, y\) by\textsuperscript{96}

\[
\theta(x, y) = \frac{\lambda_p}{2\pi} \nabla \Phi(x, y).
\]  

(3.8)

With an assumed gradient only in the \(y\)-direction, Eq. 3.7 shows that by measuring the components of the local refraction of a probe ray, the electron density profile can be determined from

\[
n_e(r) = -\frac{2}{\pi} n_{cr} \int_r^R \theta(y) \frac{1}{\sqrt{y^2 - r^2}} dy.
\]  

(3.9)

A change of variable to \(s = \sqrt{x^2 - R^2}\) was made to avoid the singularity at \(x = R\).\textsuperscript{97}

The relation in Eq. 3.9 shows that by measuring the components of the local refraction of a probe ray, the electron density profile can be determined. This formula holds for
both cylindrical symmetric and spherical symmetric plasma plumes.

\[ n_e(r) = -\frac{\lambda_p n_{cr}}{\pi^2} \int_0^{R^2-r^2} d\Phi \frac{1}{s} ds \]  \hspace{1cm} (3.10)

\[ n_e(r) = -\frac{2}{\pi} n_e r \int_0^{\sqrt{R^2-r^2}} \frac{\theta(s)}{\sqrt{s^2+r^2}} ds. \]  \hspace{1cm} (3.11)

A in using Eq. 3.11 is that the refracted angle measure used in the Abel inversion is the component of the refracted beam in the direction of the Abel inversion (\( \theta \cdot ds \)).

Measurements of refractive index of any medium are most often made by an interferometer. An interferometer is an optical device in which the coherent addition of two waves is measured. The intensity modulation that is observed is related to the relation of phase between the two overlapped beams. Using Eq. 3.4 the electron density can be directly measured from the modulation in intensity of the interferometer. For a large plasma like the one considered in this study, the electron density varies over the field of view from \( 10^{19} \) to \( 10^{21} \) cm\(^{-3} \) and the integrated path for the plasma is \( \sim 1 \) mm. The phase shift over the field of view from a plasma of this size is 1000 radians. On the interferometer, a fringe occurs every \( 2\pi \) radians of phase. Over the field of view, this corresponds to 150 fringes per mm, or 6 \( \mu \)m per fringe shift, which is too closely spaced to be resolved by the \( 4\omega \) imaging system. The spatial resolution is too poor to resolve the individual fringe shifts, making interferometry impractical for large plasmas. Figure 3.7 shows the simulated interferogram of a long scale length plasma with an embedded channel. The region most interesting to fast ignition is totally obscured by the blurring of fringes in the image. In the next section, a method to quantitatively assess the phase shift due to the large gradient in refractive index is introduced.
3.2. OPTICAL PROBING OF PLASMA

Figure 3.7: Interferometry poses a challenge for long scale length plasmas. A simulated interferogram of a channel inset in a long scale length plasma viewed from the side is shown in the top image. The simulated interferogram is calculated using $I = \frac{1}{2}[1 + \cos(2\pi fy + \Phi(x, y))]$, where $I$ is the intensity of the image, $f_y$ is the spatial frequency of the unperturbed fringes and $\Phi$ is the phase imparted onto the probe beam (bottom). The spatial resolution of the optical system is taken into account by convolving the simulated interferogram with a 3.6 $\mu$m FWHM gaussian to obtain the image. The individual fringes are too closely spaced together in the area surrounding the tip and end of the channel. The bottom image shows the phase of the beam used to generate the interferogram.
3.3 Diagnostic Set-up and calibration

To get around the interferometry problem, a new method of optical probing was introduced. There are other ways to obtain information on the spatial variation of the refractive index across a probing beam by taking advantage of the refraction angle. The deviation of the different regions of the beam may be quantitatively assessed and related to the phase shift of the probe beam and ultimately the refractive index. In this technique no reference beam is required; the variations in intensity in the image arise by the selective filtering of refracted light. There is a distinction between this method and interferometry, the direct phase information is not accessed but rather the gradient of the phase information is obtained (Eq. 3.8).

AFR is a part of a suite of diagnostics of the fourth-harmonic probe system\(^1\). A simplified optical schematic of the collection and diagnostic system is shown in Figure 3.8. In this model, the optical system is condensed into one lens with an effective focal length \(f_{eff}\). The object is illuminated by a point source located at 85.4 mm away from the object.
The diagnostic uses a filter [Figure 3.9] in the focal plane of the collection optics with alternating transparent and opaque rings to block out some of the light in the focal plane. The diagnostic relies on the correspondence between the radial location in the focal plane with the angle of a ray in the object plane. This relation is exploited by the filter that selectively blocks ranges of refraction angles, leading to shadows in the image (see Fig. 3.8). The rays are blocked or allowed shown by the red dotted and blue solid lines. Rays that are blocked lead to shadows in the image. For a single-lens imaging system and a collimated probe beam, it can be shown that a ray refracted at the object passes through the back focal plane at a distance \( r \) from the optical axis given by,

\[
r = f \theta
\]  

where \( f \) is the focal length of the lens and \( \theta \) is the angle the probing ray with respect to the optical axis before the optical collector. This relation holds regardless of the ray’s spatial location in the object plane and assumes the angles are small, \( \sin(\theta) \approx \theta \). In the more complex case of probing rays emanating from a point source as illustrated in Figure 3.8, the distance \( r \) is,

\[
r = f \frac{(M + 1)(d_s + d_o)}{d_s(M + 1) + d_o} \theta = f_{eff} \theta
\]  

where \( M \) is the magnification of the imaging system, \( d_o \) is the distance from the object to the collector, and \( d_s \) is the distance from the point source to the object plane. In the case of the collimated beam, the source distance \( d_s \) approaches infinity and \( f_{eff} \)
approaches $f$. The proportionality between the spatial location in the focal plane and the additional refraction added by the object plane allows for rays to be selectively analyzed based on their refraction angle, $\theta$.

The AFR diagnostic was calibrated by determining the effective focal length of the optical system, $f_{\text{eff}}$, in Figure 3.8. The focus of the collector is relayed $\sim$ 5 meters to the optical table, passing through fifty-eight optical surfaces before arriving at the filtering plane. The diagnostic was calibrated in-situ with a plano-concave lens with a focal length of -20 mm that was placed at the object plane of the collector. The positioning of the lens at the object plane was verified on the OMEGA EP target viewing system with an accuracy of $\pm 20 \mu$m. The lens imparts a well characterized refraction angle as a function of position onto the probing laser. A schematic of the calibration configuration is shown in Figure 3.10. By measuring the radial distance of each stripe in the image, the refracted angle can be related to the position in the filter plane. Re-
fracted light is either blocked by the opaque rings or transmitted to the image, forming a concentric ring pattern.

To infer $f_{\text{eff}}$, we introduce a controlled source of angular divergence $\theta$ into the system, with

$$r_{\text{filter}} = f_{\text{eff}} \theta.$$  

(3.14)

The calibration lens (Figure 3.10) imparts a divergence to the source beam by

$$\theta = \frac{r_{\text{object}}}{f_{\text{lens}}}$$  

(3.15)

where $r_{\text{object}}$ is the spatial location in the object plane and $f_{\text{lens}}$ is the focal length of the lens. Substituting Eq. 3.15 into Eq. 3.14 yields

$$f_{\text{eff}} = \frac{r_{\text{filter}}}{r_{\text{object}}} f_{\text{lens}}$$  

(3.16)

where all the parameters on the right hand side of Eq. 3.16 are measured. Figure 3.11 shows the results of the calibration run with a calibration lens of $f_{\text{lens}}$ of -20 mm. The value for the effective focal length $f_{\text{eff}}$ is 156.3 ± 0.5 mm.

Figure 3.12(a) shows the angular filter with the regions illuminated by a calibration lens. The orange circle shows the area of the filter illuminated by a spherical lens. A cylindrical lens forms a line focus on the filter, causing angular filtering in one direction only. Figure 3.12(b) is the image generated by the spherical lens illuminating angular filter. Figure 3.12(c) shows the calibration using a cylindrical lens of the same power as
in case (b), where the probe rays are refracted in the $y$-direction only. The focal spot in the filtering plane is elongated in one direction only, which illuminates the filter along one axis only. When the filtered signal is imaged, the filtered regions form the dark stripes in the image. The stripes are contours of constant refracted angle by the probe.

### 3.3.1 Test results

Figure 3.13 shows the image of a plasma obtained by AFR. The plasma was driven by 2 kJ UV energy in a 1-ns square pulse incident on a 125 $\mu$m thick plastic foil. The plasma was probed at 1.0 ns at the end of the UV drive. The contours of total refraction show the shape of the plasma plume as it expands from the surface of the target. The structure seen inside of the transmitted bands is a result of diffraction from the sharp edged angular filter. The effect of the diffraction can add additional uncertainty to the specific location of the contours, so the centers of the bands are used in the analysis. Data similar to Figure 3.13 has been previously obtained using a modified Schlieren setup\textsuperscript{99}. AFR expands on that technique by introducing multiple bands where as the work of Seka \textit{et al.}\textsuperscript{99} generated only a single band, limiting the region where data may be taken.
Figure 3.11: (a) The image formed by placing a negatively spherical lens in the object plane of the AFR system. (b) A radially averaged lineout (blue curve) of the image in (a). The peaks are fit with gaussian functions (red) and their centers compared with the center of the bands in the filter (green points). (c) The slope of the relation between the centers of the bands in the image and the filter bands gives $f_{eff}$ for the AFR system. The blue points are the measured centers and the red line is a fit through the points. The experimentally found $f_{eff}$ is $156.3 \pm 0.5$ mm.
Figure 3.12: (a) A photograph of an angular filter that consists of a central obscuration with a 500 μm diameter surrounded by concentric rings with a spatial period of 2 mm. (b) The calibration image of a \( f_{\text{calibration}} = -20 \text{ mm} \) focal length lens placed in the optical system. The transition from light to dark in the image corresponds to the specific light and dark circles in (a). The orange circle in (a) shows the region illuminated by the spherical calibration lens. (c) The image formed by a cylindrical lens shows that the beam is refracted in one direction only. Since this diagnostic measures the component of refracted angle only, an exact same image as (c) would be formed with a filter consisting of vertical lines and a spherical lens. This is because the filter only illuminates a small line region of the filter in this case. The green dotted line shows the region illuminated by the cylindrical lens.

A map of the integrated phase can be calculated by solving Eq. 3.8 subject to the boundary condition \( \Phi = 0 \) at the edges of the image. The bottom image of Fig. 3.13 shows the result. Eq. 3.8 is an Eikonal equation and is solved using a multi-stencil fast marching method\(^{100}\). The angle \( \theta \) at any location in the image is found by interpolating between the contours. It is instructive to note that an absolute phase shift of up to 1000 radians is observed by this diagnostic before the light is refracted outside of the collector. An interferometer would have to resolve \( \sim 150 \) fringe shifts per millimeter. This illustrates the difficulty associated with probing a highly-refractive plasma with an interferometer as it was discussed in section 3.2.

The electron density is then found by using the inferred phase map \( (\Phi(x, y)) \) from Fig. 3.13 with the help of Eq. 3.9. The result of that calculation is plotted in Figure 3.14. The shape and size of the plasma plume is consistent with a prediction from a 2-D radiation-hydrodynamic simulation with the code \( \text{DRACO}^{55} \) for the same laser condi-
Figure 3.13: (a) An AFR image obtained by probing a plasma driven by 2 kJ UV energy in a 1-ns square pulse. (b) Phase map inferred from (a).
Figure 3.14: The inferred electron density distribution (top) compares closely in size and shape to a radiation hydrodynamics simulation output (bottom).
tions that is plotted in Fig.3.14(b). Differences in the shape between the experimentally measured plasma plume profile and the simulation may be due to three dimensional effects, however the on-axis density and density scale length are in good agreement (see Fig. 4.2 in Chapter 4).

It is important to register the image with respect to the original position of the target surface when comparing with a simulation. To this end, a background image is taken immediately before the target is driven. The undriven target produces a shadow on the camera. The position of the front surface is found by locating a fiber glued to the rear surface of the target which can be seen in Figure 3.15. The fiber is located at the center of the target. The edges of the target cannot be reliably used to register the surface, as the target dimension (mm) is larger than the diagnostic depth-of-field. This causes the edges of the target to be out of focus. Before the target was inserted into the target chamber, the separation between the end of the fiber and the target surface was measured and, the position of the undriven target is accurately determined to ± 10 μm.

3.4 Fourier Optics Model of AFR

The previous section described an adequate model of AFR from a geometric optics point-of-view. This section develops a more sophisticated model of AFR based on Fourier optics that includes the effects of diffraction.

Experiments by Abbe and Porter showed that coherent image formation contains all of the spatial frequency components of the image in the back focal plane of a single-lens imaging system. Figure 3.16 illustrates the concept of their results. The Fourier spectrum (spatial frequency spectrum) of the object appears in the back focal plane. When the wavefront is propagated to an image plane, the spatial frequency components recombine to form the image. The mathematical relation between the object
Figure 3.15: The target surface is registered by measuring the location of a SiC fiber tip attached to the rear side of the target. The target surface is found by measuring the distance from the front target surface to the fiber tip before driving the target.

plane and the back focal plane (called “focal plane” in Figure 3.16) of a monochromatic wavefront is given by\(^98\)

\[
E_f(\xi, \eta) = \int_{-\infty}^{\infty} \int_{-\infty}^{\infty} E_0(x, y) \exp\left[-\frac{2\pi}{\lambda_p f_{\text{eff}}} (x\xi + y\eta)\right] dx dy,
\]

where \(E_0(x, y)\) is the electric field distribution in front of the lens in the object plane, \(\xi, \eta\) are the spatial coordinates in the focal plane, and \(\lambda_p\) is the wavelength of the light. Thus, up to a quadratic phase factor, the focal-plane amplitude is the Fourier transform of that portion of the input subtended by the projected lens aperture.

The technique of spatial filtering can be demonstrated by removing specific spatial frequency content in the Fourier plane. The periodicity of the grid creates a regular pattern in the focal plane of independent frequency components at integer multiples of
the grid’s spatial period. This effect is seen in Figure 3.17 where the image of a grid target is shown in the image plane (top) and the focal plane (bottom). All of the spatial frequency content in a given direction gives rise to the modulations in the image in that direction. Figure 3.18 shows the effect of removing all \( x \) or \( y \) directed spatial frequency components from an image of a wire grid. Fig 3.18 shows the removal of all content in one direction can totally eliminate modulation in that direction. This result can be extended to the angular filter.

The angular filter can be thought of as a selective spatial filter. In the case of a filter with concentric rings, rather than filtering \( x \) or \( y \) only, the filter blocks radial spatial frequencies that satisfy \( f_r = \sqrt{f_x^2 + f_y^2} \) where \( f_x \) and \( f_y \) are the spatial frequencies in the \( x \) and \( y \) directions respectively. The spatial frequency content in the wavefront is related to the geometric ray angles by

\[
\theta = \lambda_p f_r, \quad (3.18)
\]
Figure 3.17: (top) An image of a wire grid taken with the $4\omega$ optical probe system. (bottom) Fourier transform of (a) shows the spatial frequency content of the grid image. The points in the spatial frequency image correspond to the discrete spatial frequencies (periodicity) in the top image.
Figure 3.18: All spatial frequency content of the grid image from Fig. 3.17 is removed in the x-direction (bottom-left) and y-direction (top-left). When the image is transformed back to the image plane, all the image content in the filtered direction is lost. The resulting images (right) show that all the structure (grid edges) in the direction of the removed spatial frequency components is lost.
therefore the ‘angular filter’ serves as a spatial frequency filter when considering the
system in a Fourier optics sense.

The objects of interest to these experiments are optically thin to the probe radiation
in the region of interest. The absorption of the probe beam is neglected in this sim-
ulation. The largest effect of the plasma on the beam is the generation of a spatially
varying phase-shift, that can be traced back to the electron density by using Eq. 3.9.
Consider a plasma object with a complex amplitude transmittance

\[ t_o = P(x, y)e^{i\Phi(x, y)} \]  

(3.19)

where \( P(x, y) \) is a pupil function for the object, and \( \Phi \) describes a phase map of the
integrated optical path, given by the map in Figure 3.19. The object is illuminated
by a probing beam with a initially flat wavefront and constant amplitude inside the
beam center that transmits through a plasma object. The beam is then passed through a
focusing lens with focal length \( f_{\text{eff}} \), performing a Fourier transform on the wavefront.
At the filtering plane the intensity distribution of the beam is given by

\[ I(f_\xi, f_\eta) = F[t_o(x, y)] \times F^*[t_o(x, y)], \]  

(3.20)

where \( F[\] \) denotes the Fourier transform of the bracketed term and the physical locations
in the focal plane, \( \xi \) and \( \eta \), are given by \( \xi = f_\xi \lambda_p f_{\text{eff}} \) and \( \eta = f_\eta \lambda_p f_{\text{eff}} \) where \( f_\xi \) and
\( f_\eta \) are the corresponding spatial frequencies. The corresponding intensity for this map
is shown in Figure 3.20.

At the filter plane, a filter with amplitude transmittance, \( t_f \), is then applied to the
wavefront, yielding a new intensity distribution
Figure 3.19: Real part (top) and chord integrated phase (bottom) of the complex amplitude transmittance (Eq. 3.19) of a probe beam passing through a plasma object. The real part defined by the spatial extent of the beam, while the imaginary part is given by the phase accumulated through the plasma.
Figure 3.20: The probe beam intensity in the filter plane is shown. The majority of the signal beam is undeflected which gives rise to a bright signal at the origin. The color scale in this image is logarithmic to bring out the low level of the deflected signal in the presence of the central peak.

\[ I(\xi, \eta) = (t_f(f_\xi, f_\eta) \times F[t_o(x, y)]) \times (t_f(f_\xi, f_\eta) \times F[t_o(x, y)])^* \], \hspace{1cm} (3.21) \]

where \((\cdot)^*\) denotes the complex conjugate of the enclosed quantity. The filter has blocked specific spatial frequency ranges, that correspond to the radii of the filter. The transmittance function \(t_f\) is shown by In the example being carried out, the amplitude transmittance of the filter is chosen to match the angular filter. The angular filter that is applied to Fig. 3.20 is shown in Fig. 3.21. The wavefront is then propagated back to the image plane by carrying out an inverse Fourier transform and is given by

\[ I_i(x_i, y_i) = F^{-1}[t_f(f_\xi, f_\eta) \times F[t_o(x, y)]]F^{-1}[t_f(f_\xi, f_\eta) \times F[t_o(x, y)]]^* \], \hspace{1cm} (3.22) \]

where \(x_i, y_i\) are the spatial coordinates in the image plane. The result is shown in Figure
Thus, the effect of the angular filter can be calculated for a given phase map by two Fourier transforms.

A benefit of this method is that the transmittance function of a simulated plasma can be rapidly generated from Eq. 3.9. The predicted electron density is integrated to obtain phase. From this imaging model, it is evident where the high frequency modulation in the image arises. The knife-edges in the angular filter cause the diffraction in the image. At the time of the manufacturing of the angular filters it was not possible to get a fully apodized or "soft-edged" filter. The apodized filter would diffract less than its hard edged counterpart. A calculation with an apodized and hard edged filter shows that the image is less obscured by diffraction, at the cost of sharp edges in the image. Figure 3.23(a,b) illustrates the different transmission between a hard edged and apodized filter, and their corresponding AFR images of the plasma object are shown in Figure 3.23(c,d).
Figure 3.22: Calculated AFR image of the phase object given in Figure 3.19. The beam intensity is modulated based on the gradients in phase in the probe beam and the specific filter. The high spatial frequency structures are from the abrupt transitions in the angular filter, shown in Figure 3.21.

3.5 Summary

An optical probe for measuring laser-driven plasmas was designed and implemented on the OMEGA EP Laser System. The resolution of the imaging system was measured to be $\sim 3.6 \, \mu m$, allowing the observation of finer structures than any previous probe. The probe was used to infer the electron density of a laser produced plasma and compared to 2-D radiation hydrodynamics simulation. The probe system is modeled mathematically using a Fourier optics model to include the diffraction seen in the images. A method for minimizing diffraction is shown. The Fourier optics model allows a synthetic AFR image to be rapidly generated from the radiation hydrodynamics simulation output. This diagnostic is the basis for probing channels produced in laser driven plasma and is used to characterize both the background density plume as well as structures found inside the plume to be discussed in the next chapter.
Figure 3.23: A soft edged filter is predicted to create less diffraction in the image than the hard edged filter. In this set of images the same plasma phase object was subject to two different angular filters. (a) The transmission function $t_f$ of the hard-edged filter. (b) The transmission function of the soft-edged filter. The soft-edged filter transmission is given by $t_f(r)$ in Fig. 3.21 without the Heaviside function. (c) The simulated AFR image of a plasma with the hard-edged filter in (a). (d) The simulated AFR image of a plasma with the soft-edged filter in (b).
4.1 Introduction

This chapter describes the measurements of a plasma channel created by a multi-kilojoule high intensity laser beam in a mm-size inhomogeneous plasma. The inhomogeneous plasma was created by the irradiation of a planar plastic target with a $\sim 4 \times 10^{14}$-W/cm$^2$ UV laser pulse on the OMEGA EP Laser System$^4$. An ultrafast optical probe beam (Chapter 3) measured the channel opening time, residual density, and transverse expansion of the channel created by the intense light pressure and heating from the high intensity laser beam. These experiments demonstrate the creation of a straight, evacuated channel in an inhomogeneous plasma. The radial expansion of the channel drives a shock wave in the plasma that expands at a speed of $\sim 1\%$ of the light speed. The experimental results are consistent with the ponderomotive hole boring model$^{5,14}$, which predicts that the forward progress of the channel is driven by the light pressure. The radial expansion rate of the channel following the heating from the intense pulse was measured and the analysis shows that the channel is heated to multi-
MeV temperatures, consistent with particle-in-cell simulations that were performed by Li et al. for comparable plasma and laser conditions. This work represents the first demonstration of channeling to overcritical densities in a plasma of mm-size and extent and is important for future integrated fast ignition experiments. Increasing the interaction length, the distance over which the high intensity laser beam interacts with plasma, might have implications for other important applications such as laser wakefield acceleration and x-ray sources.

The next two chapter is organized as follows: Two experiments are discussed. The first experiment demonstrates the creation of straight, stable channels in a fast ignition-relevant plasma atmosphere. The difference in penetration extent between two pulses with different durations is investigated. The experimental results are compared to 2-D DRACO simulations to give insight into the heating and expansion of the channel. The second experiment described in the following chapter builds on the work of the first. A laser-produced channel is used to guide a second laser pulse in an inhomogeneous plasma. The results are then compared to 2-D and 3-D particle-in-cell simulations.

4.2 Motivation for the Study

The propagation of a laser beam at relativistic intensities (\(>10^{18}\) W/cm\(^2\)) through a plasma with a large density scale length is dominated by highly nonlinear interactions including ponderomotive expulsion of electrons, channeling, and the development of hosing and bifurcation instabilities. These effects are important for both fundamental aspects of relativistic laser-plasma interaction physics and applications such as fast ignition in inertial confinement fusion. The central idea of the fast-ignition concept is to first compress a frozen DT-ice capsule with a nanosecond, megajoule laser to a high areal density and then use an ultrapowerful short-pulse laser
to subsequently ignite the fuel. Since laser light at nonrelativistic intensities propagates in the plasma corona up to only the critical density \( n_c = \frac{\omega_0^2 m_e \varepsilon_0}{e^2} \) (\( \omega_0 \) is the laser angular frequency, \( m_e \) is the electron mass, \( \varepsilon_0 \) is the permittivity of free space, and \( e \) is the elementary charge), the idea is to use one high intensity pulse to form a channel through the corona and then inject the ignition pulse into the lower density plasma column as in an optical waveguide\(^{69} \) to deposit sufficient energy in the core for ignition. In contrast to the cone-in-shell concept\(^{45} \), the channeling concept has the advantage that it uses symmetric implosions and can be readily applied to cryogenic targets. Cryogenic cone-in-shell targets are technically challenging and so far have not been demonstrated. Channeling into dense plasmas relies on the laser intensity to provide sufficient ponderomotive pressure against the outflowing plasma and a long-enough laser pulse duration to sustain the channel formation. When the laser reaches densities \( > n_c \), it may continue to push forward through its ponderomotive pressure ("hole-boring") and relativistic transparency. The ponderomotive hole-boring velocity\(^{14,5} \) is found by balancing the light pressure against the pressure arising from the material stagnating against the head of the channel. For an increasing laser intensity, a higher hole-boring velocity is obtained. For a laser system of fixed energy, increased intensity comes at a price of decreased laser pulse duration or smaller spot size, and the balance between laser intensity and duration must be optimized to provide the longest channel. Channeling experiments with short\(^{43} \) and long\(^{66} \) laser pulses were performed, and it was demonstrated both in experiments\(^{65} \) and simulations\(^{18} \) that channels have a higher transmission for a trailing pulse compared to an unperturbed plasma. These experimental observations were carried out using a variety of diagnostics including interferometry\(^{65} \), self-emission\(^{60} \), and x-ray grid image refractometry\(^{61} \). Those diagnostics were limited to plasma densities below \( n_c \) [here, \( n_c \) always refers to infrared (IR)].
laser light \( (\lambda_{\text{IR}} = 1.054 \ \mu\text{m}) \). The channel region in the vicinity of \( n_c \) was unexplored because of the strong refraction of the probe radiation at high plasma densities.

The experiments here describe the observation of laser channeling in millimeter-sized inhomogeneous plasmas by measuring the extent of a channel for laser pulses with peak intensities between \( \sim 1 \times 10^{19} \) and \( \sim 4 \times 10^{19} \ \text{W/cm}^2 \). The density scale length is comparable to those obtained in high compression shots with spherical shells on OMEGA. This experiment is relevant to future integrated fast-ignition channeling experiments. To our knowledge this represents the first measurements of the channel up to \( n_c \) in a laser-driven blow off plasma of this size. The observations reported here were made possible by the use of a probe with a short wavelength \( (\lambda_p = 0.263 \ \mu\text{m}) \) and a sufficiently large solid angle of the collection optics \( (f/\# \approx 4) \). Measurements of the plasma density in the channel and in the background plasma are presented. The time for the short pulse to reach \( n_c \) was measured and compares well to simulations. The experimental results show that for a fixed laser energy a lower-intensity (but still relativistic), longer pulse propagates deeper into a long-scale-length plasma.

### 4.3 Experimental Method

#### 4.3.1 Laser system configuration

Figure 4.1(a) shows a schematic of the interaction and probing geometries. Two ultraviolet (UV) \( (\lambda_{\text{UV}} = 351 \ \text{nm}) \) laser beams smoothed by distributed phase plates (eighth-order super-Gaussian with 800-\( \mu\text{m} \) full width at half maximum)\textsuperscript{106} irradiated a 125-\( \mu\text{m} \)-thick planar plastic (CH) target to create and heat a blowoff plasma. The UV irradiation delivered 2 kJ of total energy in a 1-ns square pulse. The channeling laser pulse was an IR beam \( (\lambda_{\text{IR}} = 1054 \ \text{nm}) \) with an energy ranging from 0.75 kJ to 2.6 kJ.
An intensity comparison was performed by using 100-ps and 10-ps pulse widths. The wavefront of the channeling beam was measured and the focal-spot irradiance map was inferred for each shot. The vacuum focal spot contained 80% of the laser energy in a 25-μm spot with peak intensities of \( \sim 1 \times 10^{19} \) and \( \sim 4 \times 10^{19} \) W/cm\(^2\) for 2.6-kJ, 100-ps and 1-kJ, 10-ps pulses, respectively. The spatially averaged laser intensities are about an order of magnitude lower than the peak intensity. The focal position of the channeling beam was set to 750 μm from the original target surface, and the corresponding electron plasma density at that location was predicted to be \( n_e = 2.5 \times 10^{20} \) cm\(^{-3}\), which is close to \( n_e/4 \). It was suggested in Ref.\(^{108}\) that focusing the laser beam to \( n_e/4 \) might provide the most-favorable condition for relativistically enhanced propagation. The probe beam\(^{15}\) was a 10-ps, 263-nm laser with 10 mJ of energy (see Sec.3.1). The relative timing between the probe and channeling pulse was measured with an accuracy of better than 20 ps on each shot. The timing of the different beams is shown in Fig.4.1(b).

### 4.3.2 Probe beam configuration

The expanding blowoff plasma was measured by using optical probing, with angular filter refractometry (AFR)\(^{86}\) (Chapter 3), which measures gradients in the refractive index \( n \) of the plasma, given by \( n = \sqrt{1 - n_e/n_c} \) for a collisionless plasma. Figure 4.2(a) shows an example of such a measurement shortly before the arrival of the channeling beam. The probe light that passed through the plasma is collected and filtered in the focal plane of the collection optics where the spatial location of probing rays depend on their refraction angle. A bullseye-patterned filter with alternating transparent and opaque rings provides iso-contours of the refraction angle in the image plane. Using AFR allows one to measure the angular deviation of probe rays while preserving
Figure 4.1: (a) Schematic of the experimental setup. A plastic target is illuminated by two UV drive beams to generate a large expanding plasma plume. Following the UV drive beams, an IR channeling laser beam is injected into the plume along the density gradient. The interaction is observed by an UV optical probe pulse that is timed to arrive at a specific delay from the start of the channeling pulse. (b) The timing of each of the lasers used in the experiment.
the fine structures in the image, which is a considerable advantage over other methods. AFR has the advantage that as a refractive method it is capable of probing large plasma volumes, whereas, e.g., interferometric techniques that rely on fringe-shift measurements are severely limited by the large phase shifts accumulated along the path (See Fig. 3.7).

4.3.3 Characterization of background plasma

The electron-plasma density profile shown in Fig. 4.2(b) was inferred from the measured angular deviation $\theta$ of probing rays assuming a hemi-spherically symmetric plasma with stratified refractive index. The analysis is described in detail in Chapter 3 and in Ref. 86. The unperturbed on-axis plasma density profile in Fig. 4.2(b) is in agreement with two-dimensional (2-D) hydrodynamic simulations with the code DRACO55. The simulated electron temperature is 1.8 keV. The experimentally determined radial density scale length in the observed region varies from 200 to 320 $\mu$m with 250 $\mu$m being the average. The last contour in the collection system ($\theta \approx 8.1^\circ$) corresponds to light that is refracted through a peak density above $n_c$ ($1.4 \times 10^{21}$ cm$^{-3}$).

4.3.4 Laser timing measurements

A high-speed x-ray streak camera109,110 measured the laser timing for each shot. The laser heated plasma emits x-rays through bremmstrahlung radiation which was recorded by the streak camera. The streak camera has a photo-sensitive cathode to convert the x-ray emission to an electronic signal, which is then recorded as a function of time. The electrons traveling ballistically through the streak tube are deflected by the voltage on the plates inside the tube. The time ramp of the voltage on the plates causes the photoelectrons to follow a different path depending on the time they are generated.
Figure 4.2: (a) Measured optical probe image of the unperturbed long-scale-length plasma (45 ps prior to the arrival of the channeling beam). The original target surface is located at $y = 0$. Contours of constant refraction angle and the associated electron densities are shown. (b) On-axis lineout of the measured values of the unperturbed electron density (points) versus distance from the target surface. The solid line is a prediction from a 2-D hydrodynamics simulation.
at the photo-cathode. After the electrons are spatially dispersed according to their generation time, the electrons strike a phosphor that converts the electron signal into a light signal. The phosphor light is collected and read out with a low-noise scientific CCD array. The streak camera was used to measure the beam-to-beam timing of the UV drive to the short pulse. PJX-II is a high speed x-ray streak camera with a temporal resolution of 20 ps when operated with a 5 ns sweep speed. The streak camera is built on the Photonis P873 streak tube\textsuperscript{111} coupled with a Spectral Instruments 1000 scientific CCD camera\textsuperscript{112}. A typical image taken from PJX-II is shown in Fig. 4.3 with the temporal axis in the horizontal direction and the corresponding time history of the laser beam overlain. The correspondence allows the timing of the two laser beams for comparison to \textit{DRACO} simulations. The vertical direction is filtered by a step wedge filter of 25, 75, and 50 $\mu$m of aluminum respectively. The image shows the x-ray emission from the generation of the plasma as a slow rise in signal over $\sim$ 1 ns, and a bright flash from the channeling beam interacting with it. The bright peak is produced by the channeling pulse in the plasma and the weaker emission generated by the UV beam. The target was the same plastic foil target described in the previous section. The laser beams were timed to arrive with a delay of 500 ps from the center of the two pulses. A fit of the trace to two gaussians gives a beam to beam timing of 580$\pm$20 ps. The measurement shows the start of the laser pulse occurs 1 ns before the peak of the x-ray emission.

4.4 Results

This section describes the AFR measurements for the 10-ps, 120-TW laser pulses and 100-ps, 20-TW pulses. The longer 100 ps pulses allowed for much deeper penetration into the plasma.

Figure 4.4 shows measured channels at different probing times for 10-ps and 100-ps
4.4. RESULTS

Figure 4.3: Time-resolved x-ray emission data obtained with PJX-II. The x-ray emission is roughly proportional to the laser intensity on-target and is used to measure relative beam timing. (a) The raw data from the CCD image measures the quantity of photoelectrons created in the streak tube. The streaked direction is horizontal and proportional to time. In the vertical direction, there is a step wedge filter of 25, 75 and 50 μm thickness aluminum. (b) An averaged lineout of the red region in (a). The streak camera signal shows the two peaks, corresponding to the 1-ns and 10-ps pulses.
laser irradiation. The channel is visualized by the perturbations in the AFR contours. The contours bend as a result of strong density gradients created by the channeling pulse. The top row in Fig. 4.4(a-d) shows the results for the 10-ps pulse. All the times reported in Fig. 4.4 are timed to the rising edge of the channeling laser pulse. At 6 ps the head of the channel reached a position 450 μm from the original target surface, corresponding to 0.6 \( n_c \). The channel was observed up to 200 ps after its creation. A comparison of Fig. 4.4(c) and (d) shows that later in time, the tip of the channel retreats backward with a velocity of \( \sim 3 \times 10^7 \text{cm/s} \) away from the target surface. There is a clear difference in the channel extent between the 10-ps and 100-ps pulses. The 100-ps pulse (bottom row in Fig. 4.4(e-h)) reaches to the contour closest to the original target surface, indicating that a density \( > 1.4 \times 10^{21} \text{cm}^{-3} \) has been reached. Fig. 4.4(e) shows that the 100-ps pulse reached in 18 ps to about the same extent as the 10-ps pulse. The 100-ps pulse continues to bore through the plasma, reaching over-critical density at 65 ps. The upper contour bands in the lower-density region are smoothly shifted in space, while the contours at higher density inside the channel are highly distorted and obscured. This is likely due to sharp density modulations at the channel wall that are observed in particle-in-cell (PIC) simulations \(^{49,64}\). Bright fourth-harmonic emission of the channeling beam was measured in the vicinity of the critical surface [Fig. 4.4(e)-(h)] with the 100-ps pulse. Harmonics from the critical-density surface have been observed in experiments with high-intensity laser beams interacting with solid-density plasma\(^{113}\). No self-generated harmonic emission was observed with the 10-ps pulse, indicating that it did not reach \( n_c \).
Figure 4.4: Optical probe images for 10-ps and 100-ps channeling laser pulses at various times showing greater penetration extent for the longer pulse. The red spots in (e)-(h) are due to fourth-harmonic generation from the 100-ps, 20-TW channeling pulse reaching $n_c$. The 10-ps, 125-TW laser pulse [(a)-(d)] never reaches $n_c$ and therefore does not emit fourth-harmonic emission. Since the detector integrated over a time much longer than the duration of the probe pulse, the harmonic signal is present even in frames taken before the channel reaches the critical surface in the probe image. Time zero is defined as the start of the short pulse. The background plasma was always the same as the one shown in Fig. 4.2(a)
4.5 Discussion

The short-pulse-laser intensity distribution in the experimental laser focus was not diffraction limited and contained a certain amount of spatial inhomogeneity, which probably seeded filamentation instability and self-focusing in the plasma driven by ponderomotive and relativistic effects. Figure 4.5 shows the profile of the laser focal spot and the dashed curve indicates the area that contains 80% of the encircled energy of the beam. Filamentation was predicted by a split-step paraxial wave equation calculation taking the ponderomotive and relativistic effects into account and using the measured wavefront map of the channeling beam and the refractive index of the plasma (Chapter 5). The calculated beam size inside the plasma was similar to the measured channel width. In addition, simultaneous measurements of the strong electrostatic and magnetic fields inside the plasma with proton radiography concluded that filamentary structures exist at a location between 0.5 and 1 mm from the initial target surface. Three-dimensional (3-D) hydrodynamic simulations including relativistic corrections and the effect of charge separation have demonstrated that aberrated beams do not channel as effectively as diffraction-limited beams. Filamentation is sensitive to the power in the laser speckles and therefore is expected to be more severe for higher-power pulses. When sufficiently driven, this might cause beam spraying and result in the breakup of the beam so that it cannot reach a higher density. The 100-ps channeling beam had \( \sim 4 \times \) less power than the 10-ps pulse and is expected to be less affected by filamentation, which might be one of the reasons why this beam propagated deeper into the plasma.

Two- and three-dimensional PIC simulations with large (\( \sim 500-\mu m \)) plasmas studied the propagation and channeling for conditions similar to the experiment. The laser power greatly exceeds the power threshold for relativistic self-focusing and beam filamentation occurs in the early stages of the simulation. The local intensity in-
increases in the filaments and the resulting transverse ponderomotive force pushes most
of the electrons out of the filaments. The resulting space-charge force causes the ions to
follow, creating several microchannels that eventually merge together and form a single
density channel along the laser propagation axis. The simulations predict that besides
laser hosing, channel bifurcation, and self-correction, the laser front will pile up ma-
terial at the channel head that will reach densities \( n > n_c \), even though the surrounding
plasma is underdense. The simulations predict that after a short initial period (\( \sim \) ps),
when the pulse propagates with a speed close to the linear group velocity, it quickly
slows down and, after \( \sim 5 \) ps, approaches the ponderomotive hole-boring velocity \(14,5\).
The plasma density gradient rapidly steepens in front of the pulse and the laser light
essentially interacts most of the time with steep overcritical plasma \(14,18\).

The channel propagation velocity can be obtained from the experiment. The extent
of the channel is found by measuring the distance from the original target surface to the
closest point of perturbed contours in the probe image. A channel progression velocity
of \( > 3 \) \( \mu m/ps \) was obtained. The Mach angle gives another measure of the velocity of

Figure 4.5: The focal spot of the channeling beam in (a) linear scale and (b) log scale. The red dotted
circle shows the radius containing 80% of the encircled energy of the beam.
Chapter 4. Channeling of Multi-Kilojoule High Intensity Laser Beams in an Inhomogeneous Plasma

Figure 4.6: The perturbation to the background electron density propagates outward in a bow shock wave showing a Mach cone. The angle of the resulting Mach cone gives an estimate of the velocity of the forward-going channel front. Here the Mach angle $\sim 9^\circ$ gives a channel head velocity of $\sim 2 \mu$m/ps.

The supersonic advancing front in the gas$^{118}$. The velocity of the front of the channel is found by measuring the angle of the wake left behind by the channel. The Mach angle relates the front Mach number $M$ to the angle $\theta_m$ by $\sin(\theta_m) = 1/M$. The Mach angle is measured to be $\sim 9^\circ$ [in Fig.4.6, also observed in Fig.4.4(f)], which gives a channel head Mach number of $M \sim 7$ and a velocity of $2 \mu$m/ps, slightly lower than from the extent measurement.

4.5.1 Channel extent into the plasma

The measured channel extent is compared to a simple model calculation. The model is constructed by balancing the forward directed light-pressure with the dynamic pressure generated by the head of the channel pushing plasma out of the way and is based on the hole-boring concept$^5$(see Chapter 2.5). The velocity increases for increasing laser intensity, and slows as the density ahead of the channel increases. The channeling velocity can then be found as,
\[
\nu_c = \sqrt{\frac{I_L Z (1 + R)}{2 n_e m_i c}}.
\] (4.1)

The extent of the channel is then calculated by balancing the forward velocity \(\nu_c\) against the backward going blow-off velocity \(\nu_b\), which arises from the laser heated plasma expanding away from the target surface,

\[
y'(t) = \sqrt{\frac{I_L(t) Z (1 + R)}{2 n_e(y) m_i c}} - \nu_b,
\] (4.2)

where \(y(t)\) is the position of the channel in the plume. This is a non-linear ordinary differential equation that describes the position of the channel front in the plasma. The blow-off velocity was measured by observing the recession of the channel after the laser was turned off. The blow-off velocity is calculated by the hydro-simulation and is found to be in agreement with the observed recession value of \(3 \times 10^7\) cm/s. Equation 4.2 is solved using a 4th order Runge-Kutta solver subject to initial conditions and parameters

\[
y(0) = 750 \, \mu m, \quad (4.3)
\]

\[
y'(0) = 0, \quad (4.4)
\]

\[
n_e(y) = n_0 e^{\frac{y}{L_s}}, \quad (4.5)
\]

\[
I_L(t) = I_0 e^{-(\frac{t}{\tau_L})^6}. \quad (4.6)
\]

The initial density \(n_0\) is \(1 \times 10^{19} \, \text{cm}^{-3}\), \(\tau_w\) is the temporal width of the laser pulse, \(L_s\) is the measured density scale-length in the plasma and \(I_0\) is the peak intensity of the laser pulse. The channeling pulse is modeled by a sixth-order super-Gaussian temporal
shape. This analysis does not take the plasma material pile-up at the channel head into account. The material pile-up is required to create a reflection front for the laser, but the work required to steepen the front is small. Figure 4.7 shows the result of this calculation for a 100-ps, $I_0 = 10^{18}$-W/cm$^2$ laser pulse from $t = 0 - 500$ ps. The laser pulse is represented by the red solid curve and the trajectory of the front given by the model is shown by the blue curve. The open points show the extent measurements from the AFR images with horizontal error bars given by the timing error in the laser system and the vertical error bars are given by the fitting of the front of the channeling pulse to the image. The calculated trajectory of the channel head is in good agreement with the observed extent of the channel.

The measured channel progression velocity is in agreement with PIC simulations showing that it approaches the hole-boring velocity, given by $v_h = c \sqrt{Z(n_e/n_c)(m_e/M_i)(2 - \eta_a)I_{18}^2\lambda_\mu^2/5.48}$ for normal incidence light$^5$. Here $c$ is the
speed of light, $Z$ the average charge state of the plasma, $M_i$ the ion mass, $m_e$ the electron mass, $\eta_a$ the laser absorption fraction, $I_{18}$ the laser intensity in units of $10^{18} W/cm^2$, and $\lambda_\mu$ the laser wavelength in units of $\mu m$. The hole-boring velocity decreases with $1/\sqrt{n_e}$ when the pulse propagates deeper into the plasma, forcing more material to pile up. For an average intensity of $I = 10^{18} W/cm^2$, fully ionized plastic, $\eta_a=1$, and $n_e/n_c$ between 1 and 2, $v_h$ is estimated with 2.3 $\mu m/ps$ and 1.6 $\mu m/ps$, respectively, which is consistent with the measurements. The 3-D simulations provide scalings for the time $\tau_c$ that is required for the channel head to reach the position of $n_c$ and the required laser energy $E_c$, which are given by $\tau_c(\text{ps}) = 1.5 \times 10^2 I_{18}^{-0.64}$ and $E_c(\text{kJ}) = 0.85 I_{18}^{0.32}$. The estimated times and energies are between $\sim 35$ ps, 1.9 kJ and $\sim 150$ ps, 0.85 kJ for $10^{19} W/cm^2$ and $10^{18} W/cm^2$, respectively, in rough agreement with the experimental values. This scaling demonstrates that the 10-ps pulse is too short and energy poor to reach $n_c$. The time it takes for the pulse to bore through one scale-length distance is $250 \mu m / (2 \mu m/ps) \sim 125$ ps. This indicates that the channeling process is driven primarily by the ponderomotive force as predicted by the 2-D and 3-D PIC simulations. The results of the full 2-D PIC simulation are shown in Fig. 4.8. The chart shows that for laser intensities around $10^{19} W/cm^2$, the channel propagates to $\sim 0.5 n_c$ in 10 ps in good agreement with our measurements. It also shows that the extent of the channel dramatically increases as the pulse duration is lengthened and holding the energy constant. This was confirmed in the experiment with the 100-ps pulse. This is illustrated by comparing the final density for intensities of $10^{19}$ and $10^{18} W/cm^2$ at the 10 and 100 ps times, respectively. The simulations were not run for long enough for the channel to reach the critical surface, however the extrapolations (solid lines in Fig.4.8) are consistent with the experimental measurement.
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4.5.2 Density inside channel

Figure 4.9 shows a radial cross section of the measured density profile in the channel. The density is calculated through an Abel inversion of the phase, which is inferred from the angular refraction in the AFR image. The density profile in the channel region takes on a parabolic-like shape bounded on either side by walls with a density higher than the background density. For the specific image shown here, the density in the channel has been reduced from $2 \times 10^{20}$ cm$^{-3}$ to $8 \times 10^{19}$ cm$^{-3}$—a reduction of $\sim 60\%$. The measured density in the channel is about a factor of 2 higher than that predicted by PIC simulations for similar conditions$^{49}$. This image was taken at 65 ps into a 100-ps pulse, so the density may still be decreasing. At later times, contours become impossible to identify as the density gradient at the channel wall steepens.

4.5.3 Shadowgraphy

Imaging the near field of the $4\omega$ probe beam after passing through the plasma provides a shadowgraph$^{119}$ of the interaction. Shadowgraphy is a simple optical method
4.5. DISCUSSION

Figure 4.9: Density profile of a channel created by a 100-ps pulse and probed at 65 ps into the pulse along the x axis at y = 0.8 mm. The density on axis is approximately reduced by 60% with respect to the density of the unperturbed plasma at the same location (n_e = 2 \times 10^{20} \text{ cm}^{-3}). Error bars are calculated from the uncertainty in the measurement of \theta.

where variations in the refractive index of the plasma are visible in an image due to the bent paths of rays passing through. A simplified setup is shown in Fig. 4.10(a). For the specific case of the optical probe, light is either transmitted if its refraction angle is less than half of the collection f/# of the optics, or rejected if the refraction angle is larger than half of the f/# of the collection. Figure 4.10(b) shows a shadograph image of a bullet fired in air\(^{17}\). The bullet in the image above is moving so fast, the air in front of it cannot get out of the way quickly enough. It starts to pile up in front of the bullet, forming an area of compressed air—a shock wave. The abrupt change in density from the shock wave acts as a refractive index jump causing the bright line in the image. The resolution of the system is much higher in this case due to the unrestricted pupil of the optical system compared to AFR. In AFR, half of the spatial frequency content is rejected at the filter, which leads to a wider point spread function for the optical system. Figure 4.11 shows the comparison between the point spread function (PSF) of the
optical system with and without the AFR filter. The presence of the additional lobes in the PSF shown in blue degrades the sharpness of the image.

4.5.4 Temperature in the channel

The radial extent of the channel was measured as a function of probing time. Here the radial extent refers to half of the distance between the bright edges at $y = 0.75$ mm [Figure 4.12] of the channel seen in the shadowgraph. Tracking the distance from edge to edge, the expansion velocity is found. The bright edge in the image is formed because material piles-up at the channel wall which increases the density (see Fig. 4.9). Figure 4.13 shows the radial expansion at the position of the laser focus (750 $\mu$m in front of the target surface). A self-similar cylindrical model is used to explain the expansion of the channel as a function of time\textsuperscript{40,120}. The self-similar Sedov cylindrical blast wave solution\textsuperscript{121} is given by

$$ R(t) = \left( \frac{E_{\text{th}}}{\rho_0} \right)^{\frac{1}{2}} t^{\frac{1}{2}} $$(4.7)
where \( R(t) \) is the radius of the cylindrical shock wave, \( E_{\text{th}} \) is the thermal deposition per unit length, \( \rho_0 \) is the mass density of the unperturbed plasma and \( t \) is the expansion time. The mass density \( \rho_0 \) is inferred from the electron density by \( \rho_0 = n_e m_i / Z \). The mass density at the laser focus is 1 kg/m\(^3\). The data in Fig. 4.13 is fit to the model given by Eq. 4.7 in the form

\[
R(t) = \alpha (t - t_0)^0.5,
\]

and the best-fit to this model is plotted in Fig. 4.13 in red for the ensemble of 10 ps data, and in blue for the 100 ps data. The fitting parameter for the 100 ps expansion data is \( \alpha = 19.7 \pm 1.8 \). The fitting parameter for the 10 ps expansion data is \( \alpha = 22.3 \pm 5.7 \). The standard deviation of the fit for the 10 ps pulse is higher due to the smaller number of data points.
Figure 4.12: Shadograph images for the 100-ps channeling pulse at various probing times. The channel boundary radially expands forming a shockwave into the unperturbed plasma material. Time zero is defined at the start (50 ps from the center) of the 100-ps pulse. The bright edge highlights the channel wall and occurs due to a local refractive index variation over the width of the shockwave due to the density jump. The wave progresses radially outward and slows according to the formula of Eq. 4.7.
of data points in the set. The thermal energy deposited per unit length of the channel is found from the fitting parameter $\alpha$,

$$E_{\text{th}} = (\alpha \rho_0)^4. \quad (4.9)$$

The two fitting parameters for the 10 and 100 ps pulses result in a thermal energy deposition per unit length of $234 \pm 150$ and $150 \pm 61$ J/mm, respectively. The AFR images show that the channel is $\sim 1$ mm in axial extent, thus the estimated deposited energy into the expanding blast wave is $\sim 5$-10% of the incident laser energy. The remaining energy is most likely reflected, consistent with the observations from Sec. 4.5.1 which assumed that most of the laser energy is reflected from the channel front.

The plasma in the channel is heated by the intense laser pulse. Particle-in-cell simulations show that the plasma is heated in the first few ps of the laser pulse. We use this result to assume that the channel is heated instantaneously. The outgoing velocity is related to the temperature of the plasma inside the channel of a freely expanding cylinder, assuming that the expansion is driven solely by the heated plasma inside the channel. The pressure of the surrounding blow-off plasma is negligible compared to the pressure inside the channel, and the shock speed is related to the temperature inside the channel by

$$C_s = \sqrt{\frac{\gamma_e Z k_B T_c}{m_i}} \quad (4.10)$$

where $\gamma_e$ is related to the ratio of specific heats and $\sqrt{\gamma_e} = 1$ in the case where the electrons behave isothermally, $T_c$ is the channel temperature and $k_B$ is Boltzmann’s
Figure 4.13: The radius of the cylindrical shock wave as a function of time. The uncertainty in position is based on the thickness of the bright edge in the shadograph and the uncertainty in time is based on the on-shot timing measurement ($\pm 20$ ps). The red symbols correspond to measurements of a channel bored with 10-ps laser pulse, while the blue symbols indicate the expansion for the 100-ps pulse. The images in (b),(e) and (f) were taken with the near-field camera with a resolution of $\sim 20 \mu$m.
constant. Figure 4.14 shows the temperature of a channel cooling by expansion work as a function of time. This is calculated by taking the derivative of channel expansion model of Eq. 4.8 and substituting into 4.10 and solving for $T_c$. The expansion may not be solely due to expansion work, radiation may contribute to the cooling of the channel as well. A radiation-hydrodynamics simulation was carried out to estimate the behavior of the channel. Three simulations were carried out where a cylinder of heated plasma was placed into the background ablation plasma. Figure 4.15 shows a temperature map of the hot channel embedded in the plasma. The background plasma conditions are based on the radiation hydrodynamics simulation of Fig. 3.13 in Chapter 3. The channel is assumed with initial temperatures of 100, 800 and 3000 keV. The temperature evolution at the center of the channel as a function of time is plotted (squares) alongside the cylindrical model in Fig. 4.14.

The data suggests that the self-similar expansion solution gives an acceptable estimate of the initial temperature in the channel and that the cooling is primarily driven by expansion and not radiative cooling. The temperature in the channel is initially close to 1000 keV and cools to $\sim$ 20 keV over the first 100 ps. Particle-in-cell simulations predict that the channel reaches $> 3000$ keV in the first few picoseconds of the channeling beam. Once the plasma inside the channel is heated to $> 1000$ keV the electrons left in the channel are no longer influenced by the laser. Long before this, collisional absorption of the beam dropped significantly at a considerably lower temperature, decoupling the ions from the electrons and by extension, the laser field. The temperature of the electrons continues to increase through parametric processes between the laser field and the electron fluid. The electron temperature increases until a certain point where the ponderomotive force felt by the electrons which is the coupling between the laser field and the electron fluid vanishes. At this point the fast electrons are fully decoupled from the effects of the laser. This is a self-limiting process, the channel temperature
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Figure 4.14: Inferred temperatures from the fitting of the channel expansion channel is shown by the solid red and blue lines. Three DRACO radiation-hydrodynamics simulations were run to compare to the expansion in the channel. The three temperatures are 100 (yellow), 800 (green) and 3000 (blue) keV. The channel expansion appears to follow an expansion assuming an initial plasma temperature of $\sim 1000$ keV.

Figure 4.15: This image shows the initial temperature profile of the 2-D cylindrical symmetric channel expansion simulation. The temperature of the plasma was initially set to 100, 800 or 3000 keV and allowed to expand freely into the ablated plume.
4.6 Summary and Conclusions

Experiments have been performed to investigate the transport of high-intensity (> $10^{18}$ W/cm$^2$) laser light through a millimeter-sized inhomogeneous kJ laser-produced plasma up to overcritical density. The high-intensity light evacuates a conical-shaped cavity with a radial parabolic density profile that is observed using a novel optical probing technique—angular filter refractometry. The experiments showed that 100-ps
infrared pulses with a peak intensity of up to $\sim 1 \times 10^{19}$ W/cm$^2$ produced a channel to plasma densities beyond critical, while 10-ps pulses with the same energy but higher intensity did not propagate as far. The plasma cavity forms in less than 100 ps, using a 20-TW laser pulse, and advances at a velocity of $\sim 3 \mu$m/ps, consistent with a ponderomotive hole-boring model. Detailed images of the channels produced by the intense light pressure of a multi-kilojoule laser beam in a plasma were taken with a 263-nm optical probe. The evolution of the channel walls is quantitatively consistent with heating the plasma electrons inside the channel to MeV temperatures. The development of a radial shock wave and the longitudinal channel formation is fully established within a 100 ps time scale. After the channeling pulse ceases, the plasma in the channel rapidly cools because of radial expansion. The expansion time and cooling are consistent with 2-D axisymmetric DRACO calculations. The duration and temperature of the channel are encouraging for applications requiring a transmitting plasma waveguide, such as for fast ignition. A benefit of the strong heating in the channel is that the plasma becomes more transmissive. Even if the density is not reduced, the high temperature decreases the attenuation of laser light in the channel, allowing for the efficient transmission of a subsequent high intensity pulse. A follow-up set of experiments were conducted to attempt to transmit a 10-ps pulse through a preformed, heated channel created by a 100-ps pulse. The next chapter discusses the results of this experiment.
Chapter 5

Guiding of a high-intensity laser pulse with a preformed laser produced channel

The ability to propagate a short pulse, ultraintense laser pulse over a long distance in a plasma atmosphere is crucial for a number of applications in the plasma physics community such as compact laser wakefield accelerators, x-ray laser schemes, and fast ignition of laser compressed fusion targets. In this chapter we demonstrate the efficient energy transport of an ultra-intense laser pulse through a non-uniform, inhomogeneous low density plasma expected in the corona of a fast ignition target. The transmission of the ultra-intense pulse is estimated by observing the depth the pulse reaches into the inhomogeneous plasma with and without a plasma channel guiding structure. Complementary studies by Fuchs et al. have observed the quenching of forward-going laser instabilities of the intense pulse, propagating in a similar plasma atmosphere. The results show that a high-intensity pulse (> 10^{19} W/cm^2) can be directed into a preformed channel and deposit it’s energy deeper into plasma than without the preformed channel.
5.1 Motivation for study

A path for a second high-intensity laser pulse may be provided by a preformed, heated channel in the inhomogeneous coronal plasma. Channels have been shown to be effective waveguides for intense laser pulses\textsuperscript{122}. The previous chapter demonstrated the technique of generating a plasma channel, which acts as an “optical fiber” embedded in a plasma through a density modulation in the plasma. The plasma refractive index channel was produced through the hydrodynamic evolution of a high-power laser focused into plasma. The focused beam heats and ejects particles from the focal region, driving a powerful radial shock wave in the electron density, leaving a low density region on axis. The dependence of refractive index on plasma density means this cavity structure behaves as an optical fiber, with a region of higher refractive embedded in a lower refractive index medium. Additionally, the heated channel has been shown to minimize laser-plasma instabilities\textsuperscript{123}.

In this series of experiments, we show efficient transport of energy through the coronal plasma using a preformed channel established in the inhomogeneous coronal plasma of a laser-ablated target. The transmission of the second pulse is nearly the same as in vacuum. Laser-plasma instabilities may be quenched by the establishment of a heated preformed plasma channel.

5.2 Experimental method

The laser and target parameters for the colinear propagation experiment were similar to those of the channeling experiment described in the previous chapter. The underdense plasma was the same as the one generated in Fig\textit{4.2}. The beam configuration is identical to Fig\textit{4.1}, with the addition of a second pulse propagated along the same axis.
as the channeling beam. Figure 5.1 shows the laser configuration and the laser timing of the four laser pulses. Five different laser beams were used in this experiment. The plasma was characterized with AFR and shadowgraphy to image the interaction. Two ultraviolet (UV) ($\lambda_{UV} = 0.351$ μm) laser beams smoothed by distributed phase plates (eighth-order super-Gaussian with 800-μm full width at half maximum) irradiated a 125-μm-thick planar plastic (CH) target to create and heat a blowoff plasma. The UV irradiation delivered 2 kJ of total energy in a 1-ns square pulse.

The third laser pulse ("channel" pulse) is injected along the density gradient to create a hollow, heated density channel (Chapter 4). The channel is observed to be evacuated to within the precision of the AFR diagnostic. The channeling pulse ($\lambda_{IR} = 1054$ nm) is a 100-ps, 1 kJ pulse focused by an f/2 off-axis parabola to a 25 μm radius spot with a focused intensity of $5 \times 10^{17}$ W/cm$^2$ average intensity 750 μm in front of the original target surface onto a density of $2.5 \times 10^{20}$ cm$^{-3}$ measured by AFR. The fourth pulse ($\lambda_{IR} = 1054$ nm), the intense interaction pulse, is sent into the shock-prepared plasma channel. The pulse is 10-ps, 1 kJ, and focused to 25 μm, with an average intensity of $5 \times 10^{18}$ W/cm$^2$ and a peak intensity of $2 \times 10^{19}$ W/cm$^2$. It is also focused 750 μm in front of the initial target. The Rayleigh length of the pulse in vacuum is $\sim 60$ μm (See Fig. 5.9), which is considerably shorter than the plasma and channel extent. The channeling and interaction pulses are delayed by $\sim 35$ ps between the centers of the two pulses. This means the interaction pulse is injected into the channel well after the channel is established (Chapter 4) and ensures that the interaction pulse always arrives after the channeling pulse has begun to create the channel (timing jitter 20 ps). Because both pulses are focused using the same beam path and parabola, they are co-aligned to the same optical axis.

The timing of the two pulses was measured using an x-ray streak camera. Figure 5.2 shows a streak camera trace showing the two co-propagated IR pulses. Prior to
Figure 5.1: (a) Schematic of the setup for the guiding experiment. A plastic target is illuminated by two UV drive beams to generate a large expanding plasma plume. Following the UV drive beams, an IR channeling laser beam is injected into the plume along the density gradient. Following the establishment of the channel, a fourth interaction pulse is injected into the channel. The interaction is observed by an UV optical probe pulse that is timed to arrive at a specific delay from the start of the channeling pulse. (b) The timing of each of the lasers used in the experiment.
5.2. EXPERIMENTAL METHOD

Figure 5.2: Time-resolved x-ray emission data obtained with UFXRSC. The x-ray emission is roughly proportional to the laser intensity on-target and is used to measure relative beam timing. (a) The raw data from the CCD image measures the quantity of photoelectrons created in the streak tube. The streaked direction is horizontal. In the vertical direction, there is a step wedge filter of 25, 75 and 50 μm thickness aluminum. (b) An averaged boxout of the red region in (a). The streak camera signal shows the two pulses, corresponding to the 100 ps and 10 ps pulses. (c) The zoomed in region in (b) shows the separation between centers of the two pulses is ∼ 35 ps.

The experiment, the channeling and interaction pulses irradiated a 500 × 500 × 20 μm copper foil. The two pulses create copious Kα x-rays that are measured by the streak camera. The distance between the centers of the two pulses are measured to be ∼ 35 ps by fitting the two pulses and finding the separation of the peaks.
5.3 Results

Figure 5.3 shows that creating a channel significantly increases the propagation length of the interaction pulse. When the preformed channel is created, the 10-ps pulse propagates further into the plasma. The second pulse is timed to arrive at the end of the channeling pulse. Section 4.4 showed that a single 10 ps pulse is stopped in the underdense corona. The 10 ps pulse from section 4.4 was $6 \times 10^{19}$ W/cm$^2$ and the simulations predict that a small amount of energy would be transported through the corona. The channel is limited in length as the inhomogeneous nature of the blow-off plasma causes the pulse to be stopped well before the initial target surface.

The interaction beam is not only better transmitted and the tip of the channel, a proxy for the spatial profile of the beam is maintained. Studies have shown that the propagation of ultra-intense lasers in plasma channels can be maintained close to the vacuum mode, owing to the fiber-optic like profile obtained from the parabolic density profile on axis in the plasma channel.

Figure 5.4 shows the AFR images of three different channels in the corona. The first image is from section 4.4 to be compared to a 100 ps, 1 kJ pulse. The third image is the effect of injecting the 10 ps pulse into the heated channel from the 100 ps pulse. It is apparent that the creation of the channel by the 100 ps pulses aids in the transmission of the 10 ps pulse.

5.4 Discussion

This result demonstrates the channel has increased transmission of the 10 ps pulse. The observed additional depth is consistent with 90-100% of the laser energy of the 10 ps pulse being transported to the end of the channel. Figure 5.5 shows the expected
Figure 5.3: The additional depth of penetration of the guided 10-ps + 100-ps pulse (top, Shot Number 17341) is contrasted with a 100-ps pulse only (bottom, Shot Number 17342). Both images are taken at the same time with respect to the channeling beam (110 ps). The pulse reaches ~ 200 μm deeper into the plasma with the interaction pulse.
Figure 5.4: (a) AFR images for the guided 10-ps + 100-ps pulse (Shot Number 17341), (b) the 100-ps pulse only (Shot Number 17342) and (c) 10-ps pulse only (Shot Number 14822).
Figure 5.5: The calculated additional distance into the plasma from the interaction pulse. The three curves are calculated by solving Eq. 4.2 and artificially lowering the laser power by 90% (red curve) and 99% (yellow curve) respectively. The results show that the head of the channel is expected to reach \( \sim 200 \mu m \) deeper for 100% transmission. The additional depth observed in the AFR image is consistent with 90-100% transmission.

In the previous chapter we showed that the resulting plasma left in the channel is heated to relativistic temperatures of \( \sim 2 \) MeV. Tzeng and Mori have shown that the remaining heated electrons in the channel effectively decouple from the laser field\(^\text{104}\). This decoupling occurs due to the increase in the relativistic mass. The electron quiver velocity in the laser field is reduced by \( 1/(1+5p_{th}^2)^{1/2} \), where \( p_{th} \) is the electron momentum spread normalized to \( mc \) for a water-bag distribution. For a non-relativistic \( T_e \), \( p_{th} \) is very large and the electron quiver motion is not reduced. For \( p_{th} \sim 10 \), the factor has dropped to 0.04. In this case, the ponderomotive force is increased by a factor of \( 10^3 \). Once the plasma channel has reached relativistic temperature, the relativistic
effects decouple the remaining plasma from the laser field. Since the ponderomotive force is the mediating force of all the parametric laser plasma instabilities, these too are quenched by the greatly diminished effect of ponderomotive force. Both ponderomotive and relativistic self-focusing are eliminated due to the decoupling of the laser field with the plasma, allowing the transmission of a second pulse in close to vacuum focusing conditions without the effect of the plasma. Simulations show that the temperature effect that is responsible for the majority of the increase in transmission of a second laser pulse. Figure 5.6 is a plot of the fractional laser power of an intense pulse propagating through a channel and without a channel obtained from PIC simulations from Ref. 17. The plot shows the fractional power of an intense pulse arriving at the critical surface as a function of time. The results show the dramatic increase in energy transport of a laser pulse through the coronal plasma is observed with a channel, shown by the red and blue curves, compared to without a preformed channel, shown by the yellow and green curves. For the pulses without a preformed channel, most of the energy is lost in the coronal plasma. The creation of a low-density channel improves the propagation, shown by the blue curve. When the temperature inside the channel is increased to 6.5 MeV in the simulation, the pulse propagates essentially with no loss through the plasma channel, showing the considerable advantage of having a high-temperature channel. The total energy transported through the plasma is < 0.001% for the green and yellow curves, ∼ 60% for the blue curve and ∼ 100% for the red curve, which is calculated by integrating each of the curves.

5.5 The Split-Step Beam Propagation Method

The effect of the channel on the wavefront of the beam may be investigated by numerically propagating the beam through an inhomogeneous plasma. In this section an
5.5. THE SPLIT-STEP BEAM PROPAGATION METHOD

Figure 5.6: The transmission of a high intensity laser pulse inside of a preformed plasma channel calculated by particle-in-cell simulations of Ref. 18. The yellow and green curves are the transmission of $10^{19}$ and $10^{20}$ W/cm$^2$ laser pulses in a FI-relevant plasma. When a channel of 0.05 $n_c$ is introduced into the plasma the transmission increases to the blue curve. If the channel is not only assumed to be evacuated but also heated to 6.5 MeV, then the transmission is shown by the red curve. The total energy transported through the plasma is $< 0.001\%$ for the green and yellow curves, $\sim 60\%$ for the blue curve and $\sim 100\%$ for the red curve.

A numerical algorithm for solving the equations governing the propagation of light in an inhomogeneous, non-linear refractive index medium is introduced. This algorithm can be used to calculated the predicted focal spot of a laser beam with a realistic near-field pattern focusing in the plasma.

The beam propagation method (BPM) is a numerical algorithm to calculate the transverse profile of a wave in an inhomogeneous medium. The paraxial wave equation can be cast in the operator form\textsuperscript{124}

$$\frac{\partial \Phi}{\partial \xi} = (L_d + L_i) \Phi,$$

where $L_d = \frac{1}{2ik_0} \nabla^2_\perp$ is a differential operator corresponding to diffraction, and $L_i = -ik_0(n - n_0)$ is the phase accumulation due to propagation through inhomogeneous
media. The solution of Eq. 5.1 for a slice of plasma material with thickness $\Delta z$ is

$$\Phi(x, y, z + \Delta z) = e^{L_d + L_i \Delta z} \Phi(x, y, z)$$

(5.2)

with the caveat that the operators do not vary with $z$ over $\Delta z$. Using the Baker-Hausdorff formula\textsuperscript{125}, the right hand side of 5.2 can be rewritten as

$$\Phi(x, y, z + \Delta z) = e^{L_d \Delta z} e^{L_i \Delta z} \Phi(x, y, z)$$

(5.3)

where $[L_d, L_i]$ can be understood as the commutation of $L_d$ and $L_i$. The terms containing the commutation of $L_d$ and $L_i$ depend on $\Delta z^2$, which we will assume is close to 0. This assumption makes the denominator of Eq. 5.3 equal to 1 and for accuracy up to order $\Delta z$ the operators are treated independently. The solution then becomes

$$\Phi(x, y, z + \Delta z) = e^{L_d \Delta z} e^{L_i \Delta z} \Phi(x, y, z).$$

(5.4)

In the absence of inhomogeneity, $L_i = 0$ and the solution can be solved using the methods of physical optics\textsuperscript{98}. One method is to convolve the initial wavefront with the free-space transfer function. In this manner, the operation of $e^{L_d \Delta z}$ is accomplished in the frequency spectral domain (k-space) by the expression

$$e^{L_d \Delta z} \Phi = F^{-1} \left[ e^{\frac{\Delta z}{2} (k_x^2 + k_y^2)} F[\Phi] \right]$$

(5.5)

where $F, F^{-1}$ represent the Fourier and inverse Fourier transform of the bracketed quantity, and $k_x, k_y$ are the $x$ and $y$ spatial frequencies. Adding in the effect of spatial inhomogeneity adds a multiplicative prefactor in the solution of $e^{-i(n(x, y) - n_0)k_0 \Delta z}$. Putting
5.6. Effect on Focal Cone

The split-step algorithm can be used to calculate the propagation of the channeling beam through the inhomogeneous plasma. The wavefront of the beam is sampled upstream and serves as the input to the split-step solver. Figure 5.8 shows the sampled
wavefront and intensity of the channeling beam near field. There are two dark vertical stripes that shield the edges of the tiled grating compressor from the energy of the beam. The wavefront is sampled using a Hartman-Shack detector\textsuperscript{107}.

A simulation domain was set up to propagate the sampled beam through a plasma. A modified version of the refractive index was used to take into account effects driven by the ponderomotive force and thermal effects in the plasma. The refractive index of the plasma is given by

\begin{equation}
\begin{aligned}
n &= \sqrt{1 - \frac{n_e(z)}{\gamma n_c}}, \\
\end{aligned}
\end{equation}

where \(n_e\) is the electron density and \(\Gamma\) is the relativistic factor given by

\begin{equation}
\begin{aligned}
\gamma &= \sqrt{1 + \left(\frac{a}{2\alpha}\right)^2},
\end{aligned}
\end{equation}

where \(a\) is the normalized vector potential (see Sec.2.6.1), \(\alpha\) represents the relativistic mass increase due to thermal spread\textsuperscript{104,127}, \(\alpha = (p + u)/(n_e mc)\), where \(p\) is the pressure, \(u\) is the internal energy density and \(n_e\) is the electron density. \(p\) and \(u\) are a function of the electron distribution function and temperature of the plasma. For the plasma under consideration here, \(\alpha\) is 1.004 for a 1 keV plasma and 8.17 for 1 MeV\textsuperscript{123}. The simulation is initialized with the wavefront of the beam in the near field and propagated through the inhomogeneous plasma with a density that follows the form

\begin{equation}
\begin{aligned}
n_e &= n e^{-z/L_s},
\end{aligned}
\end{equation}

where \(L_s = 250 \, \mu m\) was chosen to match the plasma scale length measured by AFR for this plasma. The total laser energy is conserved in this simulation (i.e. collisional
Figure 5.8: The near field intensity (top) and wavefront (bottom) of the OMEGA EP laser beam is measured on each shot by the focal spot diagnostic. The measurements are used as the input to a propagation code to calculate the focal spot of the beam in the plasma.
absorption is not included) and the plasma density is stationary. This simulation represents the earliest time in the interaction, before any channeling has occurred. Once the channel begins to form, the behavior of the plasma may change the wavefront of the beam. Four separate simulations were run to show the effects of the plasma on the channel beam’s focus. Each simulation propagates the laser field to the critical surface in the plasma and then the simulation ends. The parameters of the simulation were kept as close to the experiments as possible. The measured beam profile was used, with laser energy and intensity scaled to the experiment. The simulation window is $400 \times 400 \, \mu m$ and the step-size between planes is $10 \, \mu m$. The radius encircling 80% of the laser energy was calculated and the effective focal cone of the laser was plotted as a function of distance into the plasma. The results of the simulation are plotted in Fig. 5.9. The $r_{80}$ envelope plot shows that the 100 TW pulse (red) is unable to reach as tight of a focus as the 10-TW pulse (blue). When the plasma temperature is increased to 1 MeV ($\alpha = 8.14$), then the 100 TW pulse (green) is able to focus close to the vacuum focusing condition (black).

Figure 5.10 compares the intensity distribution of the laser at the focal plane for the 100-TW pulse without and with a channel heated to 1 MeV. In the top image, the beam has separated into individual filaments and the energy is spread into an area with a radius twice as large as when the channel is present, shown in the bottom image. The thermal effect parameter $\alpha$ effectively reduces the non-linear effect of the laser intensity on the refractive index (Eq. 5.7) allowing a higher intensity pulse to propagate without becoming unstable to filamentation in the plasma.
Figure 5.9: Comparison of focal envelopes for the beams propagating through the plasma. When a 100 TW pulse is propagated through the plasma, modulations in the wavefront are amplified by the plasma causing the beams to come to a poorer focus (red) than in vacuum (black). A 10-TW pulse is much less susceptible to the filamentation instability (blue). When the plasma temperature is increased to 1 MeV, the 100 TW pulse focuses to a smaller spot (green).
CHAPTER 5. GUIDING OF A HIGH-INTENSITY LASER PULSE WITH A PREFORMED LASER PRODUCED CHANNEL

Figure 5.10: The comparison between the focus of the 10-ps and 100-ps laser pulses having propagated through the plasma. (top) The 100 TW laser beam profile is unstable in the plasma and undergoes filamentation before reaching the focus. (bottom) When the same simulation is run in a high temperature plasma ($\alpha = 8.14$), then the nonlinear effect is reduced and the beam comes to a narrower focus.
5.7 Summary and Conclusions

The establishment of a preformed heated channel in the inhomogeneous plasma improved the propagation of a short, intense laser pulse. PIC simulations predict a nearly 100% coupling through the plasma, and experiments demonstrate that the transport of the higher intensity pulse is aided by the creation of a channel. A 100-ps pulse is preferred due to the increased penetration depth for the initial channel establishment. The transverse profile of the laser is also maintained close to the vacuum focusing condition for a heated channel.
Chapter 6

Conclusion

The behavior of ultra-intense light propagating through an inhomogenous plasma was studied in this work. Channels were created in the inhomogenous plasma, allowing the efficient propagation of a second laser pulse through the coronal plasma. To date, these experiments represent the first comprehensive studies of channel formation and guiding in fast ignition relevant plasma. Angular filter refractometry provides a new technique for probing the ultra-short dynamics of the channel formation and subsequent expansion with high spatial and temporal resolution.

An optical probe system (angular filter refractometry, AFR) for measuring laser-driven plasmas was designed and implemented on the OMEGA EP Laser System. The resolution of the imaging system was measured to be $\sim 3.6 \, \mu m$, allowing the observation of finer structures than any previous probe. The probe was used to infer the electron density of a laser produced plasma, which was compared to 2-D radiation hydrodynamics simulation. The probe system is modeled mathematically using a Fourier optics model to include the diffraction seen in the images. A method for minimizing diffraction is shown. The Fourier optics model allows a synthetic AFR image to be rapidly generated from the radiation hydrodynamics simulation output. This diagnostic is the basis for probing channels produced in laser driven plasma and is used to
characterize both the background density plume as well as structures found inside the plume. Detailed images of the channels produced by the intense light pressure of a multi-kilojoule laser beam in a plasma were taken with the 263-nm optical probe.

Experiments have been performed that investigate the transport of high-intensity ($>10^{18}$-W/cm$^2$) laser light through a millimeter-sized inhomogenous kJ laser-produced plasma up to overcritical density. It was found that the high-intensity light evacuates a conical-shaped cavity with a radial parabolic density profile. The experiments showed that 100-ps infrared pulses with a peak intensity of $\sim 1 \times 10^{19}$ W/cm$^2$ produced a channel to plasma densities beyond critical, while 10-ps pulses with the same energy but higher intensity did not propagate as far. The plasma cavity forms in less than 100 ps, using a 20-TW laser pulse, and advances at a velocity of $\sim 2-3$ μm/ps, consistent with a ponderomotive hole-boring model. The evolution of the channel walls is quantitatively consistent with heating the plasma electrons inside the channel to multi-MeV temperatures. The development of a radial shock wave and the longitudinal channel formation is fully established within a 100 ps time scale. After the channeling pulse ceases the plasma in the channel cools rapidly and is dominated by radial expansion. The expansion time and cooling are consistent with 2-D axisymmetric DRACO calculations. The long duration of the channel and the high plasma temperature inside the channel are advantageous for applications requiring a transmitting plasma waveguide. A benefit of the strong heating in the channel is that the plasma becomes more transmissive for a subsequent pulse. The high temperature decreases the attenuation of laser light in the channel, allowing for the efficient transmission of a subsequent high intensity pulse. An experiment was conducted to inject a 10-ps pulse through a preformed, heated channel created by a 100-ps pulse.

The establishment of a preformed heated channel in the inhomogeneous plasma improved the propagation of a short, intense laser pulse. PIC simulations predict a nearly
100% coupling through the plasma, and our experiments demonstrate that the transport of the higher intensity pulse is aided by the creation of a channel. A 100-ps pulse is preferred due to the increased penetration depth for the initial channel establishment. The transverse profile of the laser is also maintained close to the vacuum focusing condition for a heated channel.

This work shows the creation of radiation-pressure driven flows in plasma. The creation of these conditions may be scaled to a regime where the dimensionless ratio of radiation pressure to thermal pressure, $2I/nTc$ is $>1$, where $n$ is the density, $T$ is the temperature, $I$ the laser intensity and $c$ the speed of light. Conditions in this regime are found in physical conditions ranging from stellar interiors and radiation generated stellar winds, to the formations of “photon bubbles” in very hot stars and accretion disks$^{6,7}$ and in high-energy astrophysical environments$^8$.

Future experiments are planned to perform a fast ignition integrated channeling experiment on OMEGA to study the physics of guiding an ignitor laser pulse through a preformed channel without re-entrant cone. Two co-propagating short pulses interact with the imploded shell, the first pulse channels to the high density region and the second pulse generates fast electrons in the direction of the highly compressed core and deposits some part of its energy in the plasma. It uses a spherical implosion and the concept is directly applicable to cryogenic targets. Detailed understanding on the laser light channeling in the under-dense plasma in the over-dense plasma were obtained in the course of this work. This knowledge is applied to the design the fast ignition integrated experiment on the OMEGA laser.
Appendix A

Abel Inversion

Two-dimensional projections (images) of a three-dimensional object measure the average value of some quantity along a path through the domain. In the case of cylindrical or spherical symmetric objects, the path is a chord through the cylinder. A central problem is the reconstruction of the object based on an available two-dimensional projection. For cylindrically and spherically symmetric objects this reconstruction is called Abel inversion. In this chapter we derive the mathematical relation between the projection and the chord-integrated measurement based on Ref.\textsuperscript{94}. Consider a cylindrical symmetric function $f(r)$ and the chord integrated measure $F(y)$:

$$F(y) = \int_{-\sqrt{R^2-y^2}}^{\sqrt{R^2-y^2}} f(r) \, dx = 2 \int_{0}^{\sqrt{R^2-y^2}} f(r) \, dx \quad (A.1)$$

Figure A.1 illustrates the geometry of the problem. The goal is to find a functional form of $f(r) = \mathbb{K}[F(y)]$, where $\mathbb{K}$ is some function operating on the chord-integrated measurements $F(y)$. We start by changing the integral along the chord $dx$ into a integral for $dr$ using $x = \sqrt{r^2 - y^2}$,

$$F(y) = 2 \int_{y}^{R} \frac{f(r) \, dr}{\sqrt{r^2 - y^2}} \quad (A.2)$$
If \( f(r) \) is assumed to decay to zero, then we may replace the integral by the infinite integral,

\[
F(y) = 2 \int_{y}^{\infty} \frac{f(r)r \, dr}{\sqrt{r^2 - y^2}}. \tag{A.3}
\]

Then Eq. A.3 can be written more compactly as

\[
F(y) = \int_{0}^{\infty} K(r,y) f(r) \, dr, \tag{A.4}
\]

where,

\[
K(r,y) = \begin{cases} 
2r(r^2 - y^2)^{-1/2} & r > y \\
0 & r \leq y 
\end{cases} \tag{A.5}
\]

The kernel \( K(r,y) \) is regarded as a function of \( r \) in which \( y \) is a parameter. The asymptote in \( K \) shifts to the right as \( y \) increases and also changes the shape of \( K \). Figure
Figure A.2: The kernel $K(r, y)$ is a function of $r$ with $y$ as a parameter. It is zero for $r < y$ and changes shape according to the parameter $y$.

A.2 shows the kernel $K(r, y)$ for two values of $y$ to illustrate the variation of $K$ with increasing $y$.

A change of variable is introduced to eliminate the change of shape of $K$, which leads to a new kernel. Setting $\xi = y^2$ and $\rho = r^2$, letting $F(y) = F_A(\xi)$, and $f(r) = F_B(\rho)$ we may rewrite A.4 as,

$$F_A(\xi) = \int_\xi^\infty \frac{F_B(\rho)d\rho}{\sqrt{\rho - \xi}} = \int_0^\infty K(\xi - \rho)F_B(\rho)\,d\rho.$$  \hspace{1cm} (A.6)

with

$$K(\xi - \rho) = \begin{cases} [-(\xi - \rho)]^{-1/2} & \xi - \rho < 0 \\ 0 & \xi - \rho \geq 0 \end{cases}.$$  \hspace{1cm} (A.7)
or for $\alpha = \xi - \rho$,

$$
K(\alpha) = \begin{cases} 
(-\alpha)^{-1/2} & \alpha < 0, \\
0 & \alpha \geq 0.
\end{cases}
$$

where the integral has now been put into the form of a convolution and $K$ is a function of the difference of $\rho$ and $\xi$ only. Using the convolution theorem, we may rewrite the equation in the Fourier domain as,

$$
\frac{1}{\sqrt{2\pi}} \tilde{F}_A = \tilde{K} \tilde{F}_B,
$$

where $\tilde{F}_A$, $\tilde{K}$ and $\tilde{F}_B$ are the Fourier transforms of $F_A$, $K$ and $F_B$ respectively. The Fourier transform of $K$ can be found to be,

$$
\tilde{K}(s) = \frac{1}{\sqrt{2\pi}} \int_{-\infty}^{\infty} e^{i2\pi \alpha s} K(\alpha) \, d\alpha \tag{A.10}
$$

$$
= \frac{1}{\sqrt{2\pi}} \int_{0}^{\infty} e^{-i2\pi \alpha' s} (\alpha')^{-1/2} \, d\alpha' \tag{A.11}
$$

and using the solution to the integral\textsuperscript{128}

$$
\int_{0}^{\infty} x^n e^{-ax} \, dx = \frac{\Gamma(n+1)}{a^{n+1}}; a > 0, n > -1 \tag{A.13}
$$
where $\Gamma$ is the gamma function. The expression for the Fourier transform for $K$ becomes

$$\tilde{K}(s) = \frac{1}{\sqrt{2\pi}} \frac{\Gamma(-\frac{1}{2} + 1)}{(i2\pi s)^{-\frac{1}{2}+1}}$$  \hspace{1cm} (A.14)$$

recognizing $\Gamma(1/2) = \sqrt{\pi}$, we arrive at the expression for the Fourier transform of $K$,

$$\tilde{K}(s) = \frac{1}{\sqrt{2\pi}} \frac{1}{\sqrt{2is}}.$$  \hspace{1cm} (A.15)$$

Returning to Eq. A.9 we find

$$\frac{1}{\sqrt{2\pi}} \tilde{F}_A(s) = \frac{1}{\sqrt{2\pi}} \frac{1}{\sqrt{2is}} \tilde{F}_B$$  \hspace{1cm} (A.16)$$

Solving for $\tilde{F}_B(s)$ and again using the definition of $\tilde{K}(s)$ we find

$$\frac{1}{\sqrt{2\pi}} \tilde{F}_B(s) = \frac{1}{\pi} \tilde{K}(s)(i2\pi s)\tilde{F}_A(s).$$  \hspace{1cm} (A.17)$$

Recognizing the left hand side is the Fourier transform of the derivative\textsuperscript{129} of $\tilde{F}_A$ we arrive at

$$\frac{1}{\sqrt{2\pi}} \tilde{F}_B(s) = \frac{-1}{\pi} \tilde{K}(s)\tilde{F}'_A(s)$$  \hspace{1cm} (A.18)$$

Transforming back from Fourier domain gives the convolution

$$F_B(\rho) = \frac{-1}{\pi} \int_{0}^{\infty} K(\rho - \xi) F'_A(\xi) \, d\xi.$$  \hspace{1cm} (A.19)$$
with

\[ K(\rho - \xi) = \begin{cases} 
(\xi - \rho)^{1/2} & \rho - \xi < 0 \\
0 & \rho - \xi \geq 0,
\end{cases} \]  

(A.20)

\[ F_B(\rho) = -\frac{1}{\pi} \int_\rho^\infty \frac{F_A'(\xi)}{\sqrt{\xi - \rho}} d\xi. \]  

(A.21)

The solution of the Abel integral equation has enabled \( F_B \) to be expressed in terms of the derivative of \( F_A \). Reverting back to the original \( f(r) \) and \( F(y) \) for the radial and chord-integrated functions gives

\[
\begin{align*}
\xi &= y^2, \\
\rho &= r^2, \\
\frac{d\xi}{dy} &= 2y, \\
\frac{d}{d\xi} F_A(\xi) &= \frac{dy}{d\xi} \frac{d}{dy} F_A(y^2) = \frac{1}{2y} F'(y), \\
f(r) &= -\frac{1}{\pi} \int_r^\infty \frac{F'(y)}{\sqrt{y^2 - r^2}} dy.
\end{align*}
\]


The quantity \( F(y) \) is called the Abel transform of \( f(r) \). Equation A.26 allows the calculation of the radial profile of the cylindrically symmetrical function from the chord-averaged measurements \( F(y) \).

In reality, measurements of \( F(y) \) are made at only a finite number of discrete \( y \)-values. Interpolation between the points is used to fill in the regions between measured points. Because the spatial derivative of \( F \) appears in the Abel inversion, the calculated \( f(r) \) is sensitive to errors in measurement of \( F(y) \). In the case of AFR,
the radial symmetric electron density \( n_e(r) \) through the plasma is the quantity we seek to measure. The phase imparted on a probing beam through the plasma is

\[ \Phi(y) = \frac{\pi}{\lambda_p n_c} \int n_e(r) \, dz, \]

analogous to \( F(y) \) in Eq. A.3. Additionally, the derivative of the phase is related to the refracted angle (Eq. 3.8) \( \theta(y) \), thus we may modify Eq. A.26 to arrive at the inversion equation for AFR,

\[
\theta(y) = \frac{\lambda_p}{2\pi} \frac{\partial \Phi}{\partial y} = \frac{\partial}{\partial y} \int n_e \, dz, \quad (A.27)
\]

\[
\frac{\partial}{\partial y} F(y) = 2n_c \theta(y), \quad (A.28)
\]

\[
n_e(r) = -\frac{2n_c}{\pi} \int_r^{\infty} \frac{\theta(y) \, dy}{\sqrt{y^2 - r^2}}. \quad (A.29)
\]
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