

ABSTRACT

Title of Dissertation: LASER WAKEFIELD ACCELERATION IN
OPTICAL FIELD-IONIZED PLASMA
WAVEGUIDES

Linus Henry Feder
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Dissertation directed by: Professor Howard Milchberg,
Department of Physics

Laser wakefield accelerators (LWFA) can support acceleration gradients orders of magnitude higher than conventional radio frequency linear accelerators. This gives them the potential to drive the next generation of accelerators for high energy physics, as well as compact accelerators for many other applications. However, in order to reach higher energies and improve electron beam quality, LWFA requires the development of plasma waveguides.

This thesis demonstrates two new all optical techniques for the creation of plasma waveguides. The first, “two-Bessel” technique uses a J_0 Bessel beam to form the core of the waveguide and a higher order J_q Bessel beam to form the cladding. In the second, “self-waveguiding” technique, the guided beam itself forms the cladding of the waveguide. Preliminary electron acceleration results using the self-guiding technique, as well electron acceleration simulations are also presented.

LASER WAKEFIELD ACCELERATION IN OPTICAL FIELD-IONIZED
PLASMA WAVEGUIDES

by

Linus Henry Feder

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Advisory Committee:

Professor Howard Milchberg, Chair
Professor Ki-Yong Kim
Professor Drew Baden
Professor Phillip Sprangle
Professor Giuliano Scarcelli

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Dedication

For Henry Eckstein and Elie Lowy

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I would like to thank Prof. Howard Milchberg for hiring a Liberal Arts major as an undergraduate research assistant 8 years ago and then putting many hours into “rewiring” my brain to be a scientist. Dr. Milchberg’s confidence in my ability and unwavering support helped me through every stumble and self-doubt in my graduate school career.

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Chapter 1: Introduction

1.1 Motivation

Particle accelerators have always been a key experimental tool for researching the most fundamental properties of the universe. “I sometimes think about the tower at Pisa as the first particle accelerator, a (nearly) vertical linear accelerator that Galileo used in his studies”, wrote Leon Lederman in his book on the search for the Higgs Boson [1]. The increasing energy of charged particle accelerators has enabled steady progress in subatomic physics over the last hundred years, enabling scientists to probe deeper and deeper into the internal structure of matter. As particle physics has pushed the limits of accelerator technology, numerous other applications have been found across science and engineering, in fields from art and archaeology [2] to advanced manufacturing and medicine [3,4]. This thesis specifically focuses on laser plasma-based techniques for electron accelerators. Aside from their major role in particle physics [5], electron accelerators have been used for wide variety of industrial applications including cross-linked polymerization for making tires [6], e-beam lithography in semiconductor manufacturing [7] and electron microscopy in biology and engineering [8]. Future higher energy electron accelerators, which may use some of the science and technology

discussed here, are seen as the next step to push our understanding of the fundamental physics of the universe even further [9].

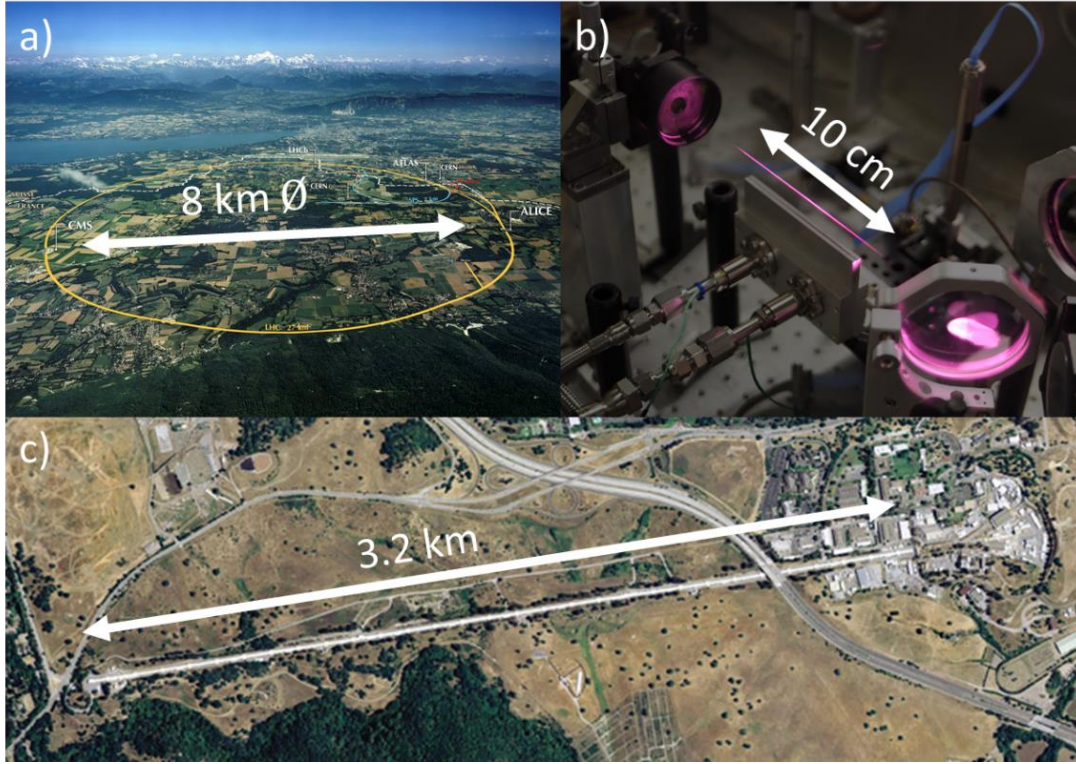


Figure 1.1 Size comparison between a) CERN AWAKE, b) UMD LWFA, c) SLAC.

In the past century, electron accelerator technology has progressed from cathode ray tubes capable of acceleration to 10s of keV to radio frequency (RF)-driven linear accelerators (linacs) reaching multiple GeV energies. RF accelerators consist of a series of resonant cavities driven by high power klystrons; the cavities can, in practical accelerators, support electric fields ~ 10 MV/m [10], with higher thresholds up to ~ 100 MV/m demonstrated in test stands [11]. This acceleration gradient is limited by the surface breakdown threshold of the cavity. Therefore, the highest energy electron accelerators are also the longest, with CERN's 27 km, 100 GeV synchrotron and

Stanford Linear Accelerator Center's 3.2 km, 50 GeV linac holding the current records. The enormous expense and engineering difficulties of building accelerators 10s or 100s of kilometers long presents a significant obstacle to future RF accelerators. Thus, accelerator physicists will need to look to alternative technologies with much higher accelerating fields for the next generation of high energy particle physics experiments.

In the years after the advent of lasers, physicists have studied the potential for using their extremely high electric fields to directly accelerate electrons [12]. However, it was shown that in vacuum the phase velocity of a focused laser beam is greater than c , therefore an electron with $v \approx c$, interacting with a laser field over an infinite distance, will not stay in phase with the accelerating field and will have no net energy gain. This result is known as the Lawson-Woodward Theorem. In 1979, Tajima and Dawson introduced plasma-based laser-driven acceleration [13], where the interaction with the plasma breaks the phase symmetry. One of the schemes proposed in [13] was laser wakefield acceleration (LWFA), where a laser pulse's pondermotive force drives a plasma wave, giving rise to an axially-directed electrostatic accelerating field generated by the charge separation between the electrons and the ions in the plasma wave. These plasma electric fields can be many orders of magnitude higher than the maximum fields in RF cavities, enabling acceleration gradients of up to 1 TV/m. After two decades of research worldwide [14], a LWFA experiment reached 1 GeV in a 3.3 cm plasma waveguide [15], a tiny fraction of the length of the SLAC accelerator. The development of LWFA has been greatly aided by the 1985 development of chirped pulse amplification (CPA) by Strickland and Mourou [16], which has enabled the construction of ultrahigh power short pulse lasers with peak powers exceeding 1 PW

(10^{15} W). This has led researchers to discuss the possibility of future laser-driven TeV electron accelerators [17].

Besides the promise of higher energies than conventional sources of electron beams, LWFA electron bunches have characteristics which open new experimental possibilities. They are ultrashort and synchronized with the laser pulse, which enables imaging of fast laser-induced processes in biology, chemistry, and physics, including other laser-plasma interactions [18].

This thesis is focused on the development and application of a new class of plasma waveguides, a key component of future LWFA experiments aiming to generate a 10 GeV accelerator stage. The remainder of Chapter 1 will introduce LWFA in more detail, discuss the need for plasma waveguides, and review the current state of the field. Chapters 2 and 3 discuss two different schemes for the generation of meter-scale, low density plasma waveguides appropriate for a 10 GeV stage, the 2-Bessel scheme and the self-waveguiding scheme. Chapter 4 reports the results of GeV-scale LWFA experiments employing the self-waveguiding scheme at the petawatt laser facility at Colorado State University. Chapter 5 discusses future LWFA experiments using our new class of plasma waveguides.

1.2 Ionization

The first step in LWFA is the formation of the plasma. Ionization by short pulses lasers can be divided into three regimes characterized by the Kelydsh parameter [19]

$$\gamma_k = \sqrt{\frac{E_i}{2U_p}} \quad (1.1)$$

where E_i is the ionization potential of an atom and $U_p = c^2 m_e a_0^2 / 4$ is the pondermotive potential of the laser and $a_0 = \sqrt{\frac{2e^2 \lambda^2 I}{\pi m_e^2 c^5}}$ is the normalized vector potential. Figure 1.2 illustrates these three regimes. For $\gamma_k \gg 1$ (left panel), multiphoton ionization dominates [20]. In this regime, ionization proceeds by the simultaneous absorption of n photons, where $n = \lceil \frac{E_i}{\hbar\omega} \rceil$ is the number of photons required to exceed the ionization potential and $\lceil \cdot \rceil$ denotes the ceiling function, rounding the number of photons up. The ionization rate is $W_n = \sigma_n I^n$ where I is the laser intensity, σ_n is the n -photon absorption cross section [20]. For $\gamma_k \approx 1$ (center panel), the laser field can be viewed as depressing the Coulomb barrier, allowing electrons to tunnel through barrier hump [21]. The ionization rate here is given by the ADK model [22]. At even higher intensities, the Coulomb field is small compared to the laser field, and the Coulomb barrier is completely suppressed. This process is called barrier suppression ionization, with the threshold given by

$$I_{BSI} = \frac{cE_i^4}{128\pi Z^2 e^2} \approx 4 \times \frac{10^9 (E_i [eV])^4 Z^{-2} W}{cm^2} \quad (1.2)$$

where Z is charge state of the daughter ion [23]. For hydrogen, the most common gas for wakefield acceleration, $I_{BSI} = 1.4 \times 10^{14} W/cm^2$, or $a_0 = .008$, far below the $a_0 > 1$ peak of a typical LWFA driver pulse. For helium-like nitrogen, a common ionization injection dopant, $I_{BSI}(N^{5+}) = 1 \times 10^{19} W/cm^2$ or $a_0 = 2.2$, which would be near the height of a LWFA drive pulse.

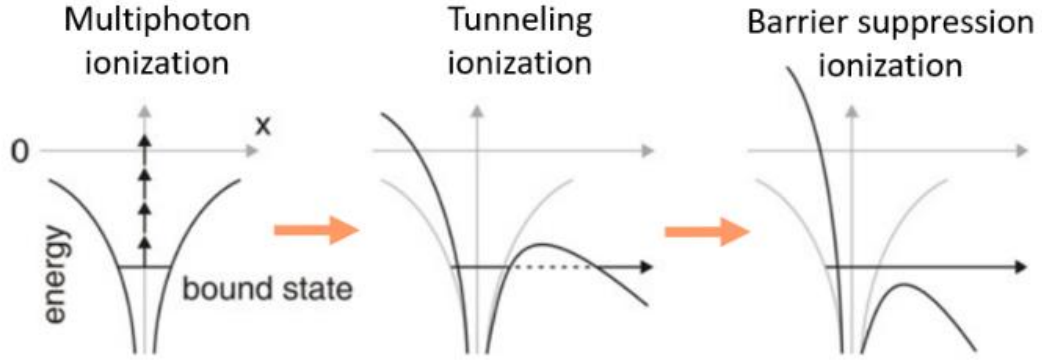


Figure 1.2 The Coulomb potential well and its deformation by an oscillating electric field in multiphoton, tunneling and barrier suppression ionization.

1.3 Plasma Wave Generation

The accelerating field in a LWFA results from the charge separation between the plasma electrons and ions in a plasma wave. The plasma wave can either be driven by a space charge of a particle beam, or, as is the case in the experiments described in this thesis, by the pondermotive force of a laser pulse. In LWFA, the negatively charged trough of the plasma wave forms a “bucket” which can trap electrons and accelerate them. This section will present a short description of the formation of the plasma wave and the coupling between laser and plasma. Beginning with the non-relativistic case, the expression for the density and electrostatic potential of the plasma waves can be derived beginning from the cold plasma equations [24,25]. We also neglect ion motion since they will remain stationary on the sub-100-fs time scale of typical laser pulses used in LWFA. For a small normalized vector potential ($|a| = e|A|/m_e c^2 \ll 1$) and small plasma density perturbations ($|\delta N_e/N_{e0}| \ll 1$), the laser ponderomotive force-driven response of the plasma density is

$$\left(\frac{\partial^2}{\partial t^2} + \omega_p^2\right) \frac{\partial N_e}{N_{e0}} = \frac{1}{2} c^2 \nabla^2 a^2 \quad (1.3)$$

where $\omega_p = (4\pi N_{e0} e^2 / m_e)^{1/2}$ is the plasma frequency and $\mathbf{F}_p = \frac{1}{2} c^2 \nabla^2 a^2$ is the laser cycle-averaged pondermotive force [14]. The solution of this equation yields sinusoidal density perturbations at the plasma frequency and a wavelength of $\lambda_p = 2\pi c / \omega_p$. The amplitude of the plasma wave depends on the laser intensity. Assuming a pulse with a Gaussian temporal profile, $a^2 \propto a_0 \exp(-(\xi/\tau)^2)$, the laser pulse is resonant with the plasma wave and drives the largest plasma wave for a given a_0 when the axial scale length of the pulse is of the same order as the plasma wavelength [14]. The electrostatic potential of the wake resulting from the plasma waves is given by

$$\left(\frac{\partial^2}{\partial t^2} + \omega_p^2\right) \phi = \frac{\omega_p^2 a^2}{2} \quad (1.4)$$

If the pulse length $\sim c\tau$ exceeds the plasma wavelength, the density gradients of the plasma wave will cause modulations in the pulse breaking it up into a train of resonant pulselets of duration $\sim \omega_p^{-1}$. This is called a self-modulated laser wakefield accelerator (SM-LWFA), and it was the first single pulse-driven scheme realized experimentally [26]. (All the experiments in this dissertation were conducted with pulses shorter than λ_p). The laser also drives transverse wakefields which have effects on the focusing of the electron beam [14] and the propagation of the laser, which will be discussed later.

For $a_0 > 1$, the electrostatic wake potential is given in the comoving frame of the laser by

$$\frac{\omega_p^2 \partial^2 \phi}{c^2 \partial \xi^2} = \frac{1 + a^2}{2(1 + \phi^2)} - \frac{1}{2} \quad (1.5)$$

$\xi = z - v_p t$ is the time in the comoving frame, since the plasma wave phase velocity $v_p \simeq v_g$, the laser group velocity [14]. Equation 1.5 also assumes the quasi-static approximation, in which the phase velocity is close to c and the laser pulse is short enough that a and ϕ change little during the transit time of the pulse [27,28]. Using Eq. 1.5 and Poisson's equation, the density perturbation $\partial N_e / N_{e0}$ and the normalized wake electrostatic field E / E_0 are computed and plotted in Fig. 1.3 (a) for $a = .5$ and in Fig 1.2(b) for $a = 2$.

$$E_0 = \frac{mc\omega_p}{e} \quad (1.6)$$

is the non-relativistic wavebreaking field (see below).

As the laser intensity increases, the plasma wave ceases to be sinusoidal as in Fig. 1.3 (a), with the electron density developing sharp peaks and the electrostatic field developing a steep sawtooth pattern in Fig. 1.3 (b). This reaches a limit – the wavebreaking limit – when the width of the density peaks becomes singular, and the electrostatic field of the wake reaches a maximum. At this point, the electron motion in a single plasma period is of order of the plasma wavelength, and the plasma wave loses coherence [29]. In the 1D, relativistic regime,

$$E_{wb} = \sqrt{2(\gamma_p - 1)} E_0 \quad (1.7)$$

where $\gamma_p = \frac{\omega}{\omega_p} = \sqrt{\frac{N_{cr}}{N_{e0}}}$ is the plasma wave's phase velocity [29] and E_0 is the non-relativistic wave breaking field. N_{cr} is the critical plasma density at which $\omega = \omega_p$ and

the laser no longer propagates in the plasma. In 3D, the wavebreaking limit is reduced by transverse wave breaking. In transverse wavebreaking, the finite spot size of the drive laser causes the electron orbits to curve towards the laser axis. At high enough laser intensities, these curved orbits cross, causing the wave to break [30].

For sufficiently high laser intensity, the electrons will be completely evacuated from the wake, forming an electron-free plasma “bubble” or cavity. This “blow-out” or “bubble regime” occurs when $2\sqrt{a_0} \gtrsim \frac{2\pi w_0}{\omega_p}$ [31], where w_0 is the Gaussian beam waist.

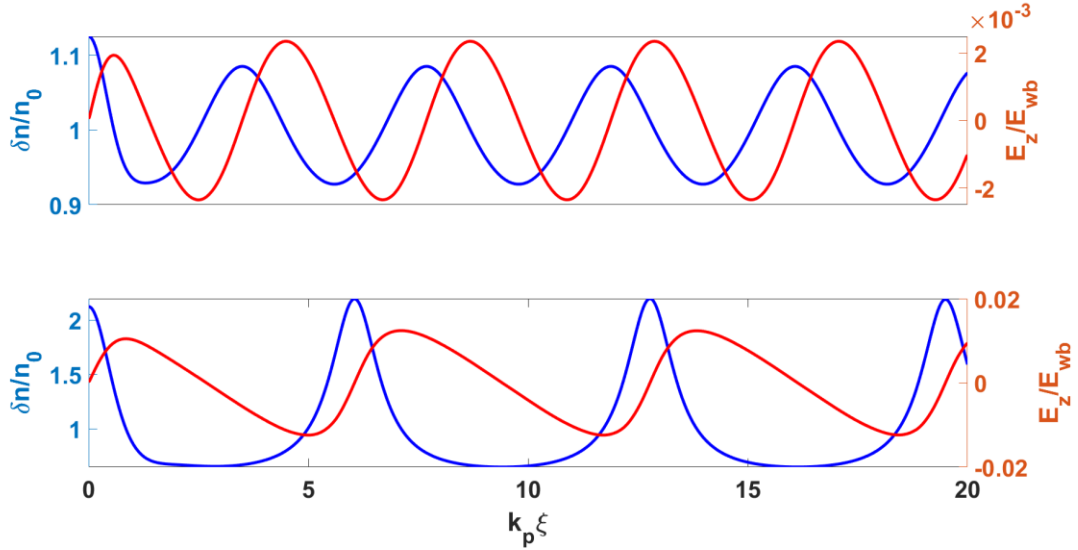


Figure 1.3. The plasma density perturbation $\delta N_e/N_{e0}$ and the normalized wake electrostatic field E/E_0 , calculated by numerically solving equation 1.5 with $a = a_0 e^{-\left(\frac{\xi}{4}\right)^2} \cos(\xi)$ for $a_0 = .5$ (top) and $a_0 = 2$ (top).

For the relativistic, 3D case, the cold fluid equations can be solved numerically using a similar fluid model to the one dimensional or low energy case [32]. However, cold fluid models break down when the particle trajectories intersect in phase space, such as in the wavebreaking process. In Chapter 4, the particle-in-cell (PIC) code FBPIC is used instead to model the acceleration process [33]. PIC codes simulate

LWFA in discrete timesteps. In each time step, the E and B fields are solved for using Maxwell's equations, then the fields are interpolated on a grid. The fields are then used to calculate the new positions and momentums of macroparticles, which represent many physical particles. Ionization rates are also calculated, and new particles deposited on the grid. The updated macroparticle positions and weights are then used to calculate the current and charge density and the loop is repeated.

1.4 Electron Injection

Although a sufficiently intense laser pulse propagating in a plasma can drive a plasma wave capable of accelerating electrons to relativistic energies, there would be no accelerated electrons without a method to trap them in the wakefield. Generally, the background electrons which form the wake do not gain sufficient longitudinal momentum in a single plasma period to stay in phase with the wakefield, so some process must be introduced to inject electrons at higher velocities. There are many different techniques for injecting electrons, each of which will affect the ultimate characteristics of the electron beam. Although the required characteristics of electrons vary depending on their use, low emittance is usually desirable. Emittance is the average spread of the electrons in position and momentum space. To minimize energy spread (the longitudinal component of emittance), the bunch will ideally be injected in a small volume and at the same phase of the wake, near the back of the first accelerating bucket behind the laser pulse, so that the entire bunch experiences the same accelerating field. Injection into several accelerating buckets or over an extended distance will increase the energy spread of the electron beam, as the electrons will be accelerated for different distances by different field amplitudes.

Conceptually, the most straightforward approach is external injection from conventional sources such as a LINAC [34], Van der Graaf accelerator [35], or RF photoinjector [36]. Such injection sources are generally poorly matched to the ultrashort wakefield acceleration bucket; aside from possible source jitter with respect to the wakefield, their electron bunches are typically much longer than a typical LWFA plasma wavelength, so their electrons are trapped in multiple accelerating buckets and experience a range of accelerating conditions, producing a wide spectrum and poor emittance. To guarantee synchronization, another laser pulse, split from the main beam, can be used as an injector. One such scheme is beat wave injection [37] in which the interference between a counterpropagating beam and the main drive beam creates a large local spike in the electric field and to inject electrons into the wake. A similar effect can be achieved by focusing a second laser for direct optical injection [38]. With both techniques, it is difficult to control the exact spatio-temporal location of injection, and laser pointing fluctuations can have a large deleterious effect.

The most straightforward and achievable injection method in practice is self-injection. If the drive pulse is intense enough to drive the wake to the wavebreaking threshold (Equation 1.7), the velocity of the oscillating background electrons reaches the phase velocity of the wake in a single plasma period, enabling them to be trapped [14,39]. Some electrons can also be trapped even below the wavebreaking threshold, because the background plasma electrons have a range of momenta based on their initial positions and momenta before interacting with the laser and the wake [40]. So, even at lower wake amplitudes, the higher momentum tail of the plasma electron momentum distribution has enough longitudinal momentum to inject into the wakefield

at the back of the plasma bubble [40]. Electrons can also be self-injected transversely through transverse wavebreaking at higher wakefield amplitudes [41]. The trapping process will continue until the trapped bunch has sufficiently high charge that it damps the electrostatic field of the wake, a process known as beam loading [14]. Without beam loading, self-injection will be continuous, assuming a uniform plasma and no depletion of the pump laser. Since electrons injected at different times and with different phase will be accelerated for different lengths, continuous injection will lead to larger energy spreads. The wavebreaking threshold depends on the plasma phase velocity (Eq. 1.7) and therefore the background electron density, so one approach to limit the region of self-injection is to longitudinally tailor the plasma density. Assuming a fixed laser intensity, a higher density injection section can reduce the wavebreaking limit enough to cause injection. This injection section is followed by an acceleration section with lower density to increase the wavebreaking limit and arrest the injection process, limiting the electron beam energy spread.

Two other injection techniques, density downramp and shock injection, rely on longitudinally tailored plasma density profiles, with their differences depending on the scale length of the density change [42]. In a plasma with a negative density gradient, the velocity of the plasma waves gradually decreases. With a sufficiently long downramp, the wake velocity will eventually decrease enough to inject electrons, albeit several accelerating buckets behind the drive laser, with the exact distance determined by initial plasma density, plasma wave amplitude and the slope of the density gradient [42]. This method, using a long density transition, is called density downramp injection and can be formed by a tailored single gas jet or the addition of an auxiliary

gas jet [43,44]. Downramp injects electrons into multiple accelerating buckets, increasing the energy spread of the electron beam and the accelerating buckets further from the drive laser have lower amplitudes, producing lower energy electron beams. Alternatively, a fast density transition, with scale length on the order of a plasma period, can be formed by a shock, which can be produced by a razor blade or other obstruction in the gas jet flow [45,46]. In this case, the longitudinal size of the plasma wave and the location of the peak accelerating field change quickly, enabling background electrons from the wave prior to the density transition to be injected [14,42].

The final injection technique discussed here, which is used in the experiments in Chapter 4, is ionization injection. In ionization injection, a higher-Z gas, commonly nitrogen or argon, is either used in the experiment as the gas target [47] or as a low density dopant of a higher density low-Z gas such as hydrogen or helium [48]. The inner shell electrons (the $1s^2$ shell of N^{5+} (helium-like nitrogen) or the $2p^6$ shell of Ar^{8+} (neon-like argon)) are not liberated until the peak intensity of the drive pulse, where they are “born” into the continuum. Since the pondermotive force is proportional to the gradient of the vector potential, the newly born electrons don’t acquire the transverse momentum of the background plasma electrons, so they can be more easily trapped by the plasma wave potential [49–51]. Like self-injection, ionization injection will be continuous for a uniform plasma and laser intensity, producing high energy spread electron beams. The longitudinal extent of injection can be limited by using a laser intensity which only exceeds the ionization threshold of the inner shell ions for a brief period [48] or by only doping a short section of the target, the approach used in Chapter

4. The amount of accelerated charge depends on the dopant concentration, with an upper limit given by beam loading.

1.5 Acceleration Scaling and Limits

The electron energy gain in a LWFA is given by

$$\Delta W = eE_z L_{acc} \quad (1.8)$$

where E_z is the axial accelerating field from the plasma wake, and L_{acc} is the acceleration distance. Limits to E_z and L_{acc} are responsible for limits on electron energy gain. The maximum E_z is limited by the laser intensity and the maximum sustainable plasma wave amplitude, given by the wavebreaking threshold, E_0 (equation 1.6). More precise values for E_z can be derived from equation 1.5, but for a resonant, 1D Gaussian pulse in the quasi-static regime, $a_0 \sim 1$ and $w_0^2/\lambda_p^2 > 2a_0^2/\sqrt{1 + \frac{a_0^2}{2}}$ [52],

$$E_{z_{max}} = \frac{.76E_0a_0^2}{2\sqrt{1 + \frac{a_0^2}{2}}} \quad (1.9)$$

At higher values of a_0 or smaller spot sizes w_0 , the accelerator enters the bubble regime, where the plasma electrons are completely evacuated from the axis, so the electrostatic potential can only be increased by increasing the size of the plasma bubble and the accelerating field is $E_{z_{max}} = E_0\sqrt{\frac{a_0}{2}}$ [31].

The other limit on acceleration is L_{acc} . Normally, a focused laser pulse is limited by diffraction to drop in peak intensity by a factor of two over the Rayleigh length, $z_R = \frac{\pi w_0^2}{\lambda}$. Diffraction can be counteracted, and the pulse peak intensity can be

maintained by using a plasma waveguide (discussed in Sec. 1.6), or by nonlinear self-guiding of the laser pulse (also discussed in Sec. 1.6). If pulse diffraction is mitigated, L_{acc} is limited either by dephasing, where the velocity of the accelerated electrons exceeds the group velocity (in plasma) of the drive pulse, or by pump depletion, where laser energy decays as it goes into the plasma wakefield. The dephasing length is defined as the length over which the electron phase slips by one half of the period of the plasma wave,

$$L_d \approx \gamma_p^2 \lambda_p \begin{cases} \frac{2}{\sqrt{\pi}} & a_0 \ll 1 \\ \frac{a_0}{2} & a_0 \gg 1 \end{cases}, \quad (1.10)$$

where $\gamma_p = \left(\frac{\omega_p^2}{\omega^2} + \frac{2c^2 \omega^2}{w_0^2} \right)^{-\frac{1}{2}} = \left(\frac{N_e}{N_{cr}} + \frac{8\pi^2 \lambda^2}{w_0^2} \right)^{-\frac{1}{2}}$ is the plasma phase velocity [14]. The second term depending on the spot size w_0 accounts for the 3D effect of diffraction on the drive laser group velocity [14]. The pump depletion length L_{pd} is found by equating the energy in the laser to the energy in the wake [53], giving

$$L_{pd} \approx \left(\frac{\lambda_p^3}{\lambda^2} \right) \begin{cases} \frac{2}{a_0^2} a_0 \ll 1 \\ \frac{a_0}{\pi} a_0 \gg 1 \end{cases} \quad (1.11)$$

For $a_0 \ll 1$, energy gain is limited by dephasing, while for $a_0 \gg 1$, it is limited by pump depletion and in the intermediate region $a_0 \sim 1$, $L_{acc} \approx L_{pd} \approx L_d$ [14]. A lower density does reduce the acceleration gradient, but assuming no diffraction and that the dephasing is mitigated by density tailoring, the energy gain $\Delta W \propto N_{e0}^{-1}$ [52], so lower densities increase energy gain. Combining the scaling of E_z and L_d or L_{pd} for $a_0 \sim 1$ gives

$$\Delta W \approx \frac{m_e c^2 a_0^2}{\sqrt{1 + \frac{a_0^2}{2} * \left(\frac{N_{e0}}{N_{cr}} + \frac{2c\omega^2}{w_0^2} \right)}} \quad (1.12)$$

For a future 1 TeV accelerator, a design of 100 stages, each of 10 GeV, has been suggested [17]. A single 10 GeV stage, based on Eq. (1.12) with $a_0 = 1.5$, requires $N_{e0} \approx 1.5 \times 10^{17} \text{ cm}^{-3}$, $w_0 \gtrsim 30 \text{ } \mu\text{m}$ and $L_{acc} < \sim 1 \text{ m}$.

1.6 Optical guiding in plasmas

Common glass optical fibers confine propagating light signals by arranging the refractive index of the ‘core’ glass (the radially central region of the fiber containing most of the light signal) to be greater than that of the ‘cladding’ glass (the region outside the core). Such solid material fibers are not suitable for guiding intense laser pulses owing to a range of severe damage scenarios, including melting, fracturing, and at high intensity flash ionization and plasma generation, followed by mechanical damage. Over a range of pulse widths, this typically limits intensity to $I < \sim 10^{10} \text{ W/cm}^2$. For LWFA experiments, where typically $I > 10^{18} \text{ W/cm}^2$, a plasma waveguide clearly must be used. The plasma, therefore, plays a dual role as the waveguide medium and as the source of the accelerating fields.

In order to provide physical insight, we can write an approximate expression, Eq. (1.13), for the index of refraction $n(r)$ as a function of radial position r for a fully ionized underdense plasma ($\frac{\omega_p^2}{\omega^2} = N_e/N_{cr} \ll 1$) in the presence of a pulse with $a_0 \ll 1$ [31]. This expression includes corrections due to (1) a preformed plasma density

profile, (2) relativistic self-focusing, and (3) a ponderomotively force-induced plasma density profile, as

$$n(r) = 1 - \frac{\omega_p^2}{2\omega^2} \left(1 + \frac{\Delta N_{channel}(r)}{N_{e0} w^2} r^2 - \frac{a_0^2(r)}{8} + \frac{\Delta N_{ponder}(r)}{N_{e0}} \right) \quad (1.13)$$

The terms denoted by (1), (2), and (3) are written in order after the “1” in the bracket of Eq. (1.13), where the “1” is the uniform plasma contribution. Examining term (1) shows that a higher density plasma has a lower index of refraction, so that a performed plasma profile suitable for guiding will have a density minimum on axis ($r = 0$). Guiding in preformed plasma profiles will be introduced at the end of this section and is the topic of the remaining chapters of this thesis, but first we will discuss self-guiding without a preformed channel. Term (3) shows that such a plasma profile can also be produced directly by the laser ponderomotive force. Term (2) is due to the relativistic effective mass increase of the quivering electrons in the intense field.

Terms (2) and (3) enable a laser with sufficient intensity to self-guide over many Rayleigh lengths, starting with a uniform gas density profile. Term (2) is the result of the change in the relativistic plasma frequency, $\omega_p = \sqrt{\frac{4\pi N_e e^2}{m_e \gamma_e}}$, due to the relativistic effective mass shift of the electrons as they are driven by the laser field, where $\gamma_e = \sqrt{1 - \frac{v_e^2}{c^2}} = \sqrt{1 - \frac{a(r)^2}{2}}$ is the Lorentz factor of the electrons. $a(r)$ is the cycle averaged normalized vector potential of the laser. Since a is peaked on axis, term (2) will act as a focusing force. Assuming a Gaussian laser spot and ignoring the pondermotive effects in terms 3 and 4, we can derive a critical power

$$P_{crit} > 17 \frac{\omega^2}{\omega_p^2} [GW] \quad (1.14)$$

above which the focusing force is stronger than diffraction [54] regardless of the initial intensity of the pulse. If the laser power exceeds the critical power, the beam will self-focus to a minimum spot size. However, at the minimum spot size, the intensity will almost always be high enough that term (3) must be included.

Term (3) in Eq. (1.13) accounts for the density depression ΔN_{ponder} created by the pondermotive force driving electrons away from the axis, where $\Delta N_{ponder} = k_p^{-2} \nabla_{\perp}^2 (1 + |a|^2)^{\frac{1}{2}}$ [55], although this formulation ignores longitudinal effects due to the non-instantaneous response of the plasma. The beam will continue to self-focus and the spot size will continue to decrease until the pondermotive force expels all the electrons, forming a bubble. For a given plasma density and laser power, the vacuum focal spot size w_0 can be chosen to match these effects, such that $2\sqrt{a_0} \sim k_p w_0$ or $a_0 \cong 2(P/P_c)^{1/3}$ [31]. In this case, the self-focused pulse will remain at a spot radius w_0 over many Rayleigh lengths. If the input beam is not focused to this so-called matched spot size, the pulse will continue to self-focus until it nears w_0 and then oscillate around that value [55].

The relative ease of implementation of self-guided LWFA compared to generating preformed plasma waveguides makes it very attractive for acceleration experiments. Some of the earliest LWFA results to achieve quasi-monoenergetic acceleration [56] and multi-GeV acceleration [57] were from self-guided LWFA. However, the power threshold $P > P_c = 17 N_{cr}/N_e [GW]$ for self-focusing makes it difficult for self-guided accelerators to operate at the low densities ($N_e < 10^{18} \text{ cm}^{-3}$) required for the highest energy gains. Figure 1.4 shows electron energy gain over a dephasing length L_d or depletion length L_{pd} , whichever is shorter, for a range of N_{e0} and a_0 using

equations 1.10-1.12. The matched spot size self-guiding condition $2\sqrt{a_0} \sim k_p w_0$ is plotted as a dotted black line. The lower density, higher energy gain phase space below the line is not accessible without a preformed plasma waveguide.

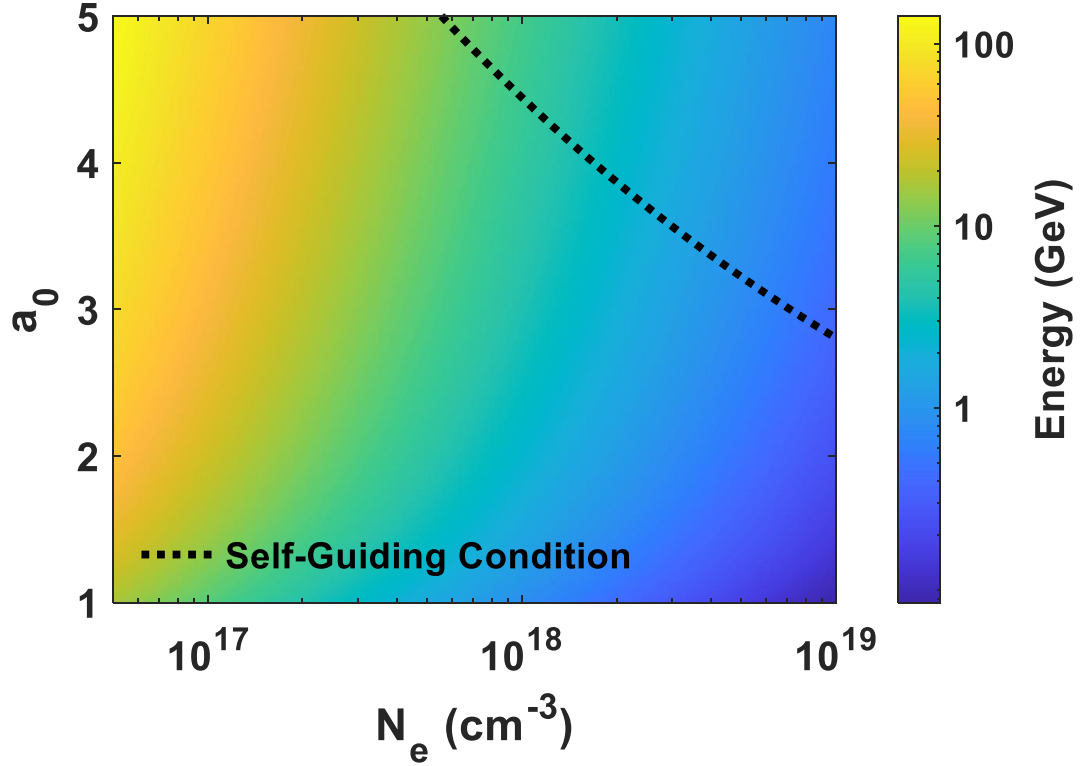


Figure 1.4. Energy gain calculated according to Eqs 1.10-1.12 for a well-matched LWFA as a function of a_0 and plasma density.

To access the phase space to the left and below the dashed line in Fig. 1.4, a plasma waveguide is required. The first plasma waveguides were created by Durfee and Milchberg in 1993 [58]. In their initial experiment, a $\lambda = 1.064 \mu m$, $\tau = 100 ps$ pulse is focused by an axicon to a J_0 Bessel line focus in a static gas fill. Through multiphoton and avalanche ionization, the pulse ionizes a column of plasma and heats it through inverse bremsstrahlung (IB) heating to $\sim 100 eV$. The hot plasma expands outward at the ion sound speed, driving a shock wave. The enhanced plasma density

of the shock walls serves as a guiding structure for a second pulse. Further improvements to this scheme followed, with the use of gas jets [59] and clusters [60], enabling more efficient laser absorption and heating. IB heated plasma waveguides were used in one of the first demonstrations of quasi-monoenergetic electron acceleration [61] and other experiments which achieved electron acceleration of hundreds of MeV [47].

An alternative approach not requiring a second laser was developed by Erlich *et al.* in 1996 [62]. In this technique, the walls of a polypropylene capillary are ionized by an electric discharge. Resistive heating of the plasma is peaked on the axis, which leads to a higher temperature and lower density at the center of the capillary than then at the walls. This expansion creates a density depression suitable for guiding. Spence *et al.* [63] refined this technique by filling the capillaries with hydrogen gas and ionizing the gas instead of the walls of the capillary, which provided a superior medium for acceleration and increased the lifetime of the capillary wall. The capillary discharge waveguide has also been combined with IB heating in order to produce a lower on-axis density [64]. That technique enabled electron acceleration to 7.8 GeV, currently a record for LWFA [65]. Recent experiments have also extended capillary discharge waveguides to 40 cm [66] and kilohertz repetition rate [67].

In the next two chapters of this thesis, two new optical field ionization (OFI)-based methods for generating plasma waveguides are presented. These methods have the advantage of significantly exceeding all earlier techniques in plasma waveguide length and, at the same time, achieving the low-axis plasma densities needed for a 10 GeV accelerator stage. These guides are also highly tunable in guided mode size, further

opening of the available phase space for LWFA. The final thesis chapter will discuss the application of these techniques to a LWFA at the 800 TW Aleph laser at Colorado State University.

Chapter 2: Two Bessel Beam Plasma Waveguides

2.1 Introduction

As was discussed in Chapter 1, the electron energy gain in LWFA scales inversely with the plasma density if the length of the acceleration is not limited by diffraction. Plasma waveguides can be used to prevent the drive laser from diffracting, but in earlier waveguide formation methods, the guided mode size is coupled to the on-axis plasma density. This chapter discusses an alternative technique which relies on two Bessel beams to directly form the waveguide with less reliance on hydrodynamics.

In capillary waveguides, this initial ionization and heating is provided by an electrical discharge and resistive heating of the plasma. The temperature deposition from resistive heating depends on electron density which couples the guided mode size to the initial gas density in the capillary, with guided mode size [68],

$$W_m[\mu m] = 1.48 \times \frac{10^5 \sqrt{R_{capillary}}}{(ZN_{gas}[cm^{-3}])^{\frac{1}{4}}} \quad (2.1)$$

After the formation of the waveguide, the on-axis plasma density depends on the initial gas density. This expression also shows that the mode size is also coupled to the size of capillary. This can be problematic experimentally, since the capillary size must be chosen based on the mode size, not necessarily large enough to prevent laser damage due to pointing fluctuations or stray light. Capillaries also suffer from erosion of their walls by the electrical discharges, which is a limit on their lifetime and repetition rate [67].

In early hydrodynamic shock waveguides, the plasma column was ionized by the laser through multiphoton and avalanche ionization. As was discussed in section

1.5, both of these processes depend on the density. In addition, to drive a strong enough shock wave to create a large enough density difference to guide, the plasma also needs to be further heated by inverse Bremsstrahlung. The heating rate of the plasma is given by

$$\frac{d\epsilon}{dt} = 2U_p \left(\frac{\omega^2 \nu_m}{\omega^2 + \nu_m^2} \right) \quad (2.2)$$

where ν_m is the momentum transfer collision rate between electrons, which is again function of electron, ion density and temperature [69]. Early hydrodynamic shock expansion waveguides operated with on-axis densities of $N_{e0} > 5 \times 10^{18} \text{cm}^{-3}$.

If clusters are used instead of regular gases, the absorption by either a picosecond scale or femtosecond scale pulses is greatly increased [60]. Clusters are solid density assemblages of atoms bound by the Van der Waals force, and clustered gas jets enable axial electron densities with $N_{e0} < 1 \times 10^{18} \text{cm}^{-3}$. However, the physics of clustering impose their own set of constraints on the gas density and species. Clustering is observed in gas jets for $200 < \Gamma^* < 1000$, where the Hagena parameter Γ^* is given by

$$\Gamma^* = \frac{k d^{.85} P_0}{T_0^{2.29}} \quad (2.3)$$

with gas temperature T_0 , pressure P_0 and nozzle diameter d [70]. k is an empirical constant with generally higher values for polarizable species. For example, $k_{argon} = 1650$ and $k_{hydrogen} = 180$, meaning much lower temperatures and higher pressures must be used to produce hydrogen clusters.

IB heating by an auxiliary laser has recently been combined with a capillary discharge waveguide to further reduce the on-axis density and give added control the

matched spot size independently of the initial gas density. A recent experiment used a 10 nanosecond heater pulse with the capillary discharge plasma waveguide to reach $N_{e0} = 3 \times 10^{17} \text{ cm}^{-3}$ for $w_0 = 60 \mu\text{m}$ [64]. Still, since the IB heating rate depends on the plasma density, this approach does eliminate the dependence of the guided mode size in capillary discharge waveguides on initial gas density.

To avoid coupling the guided mode size to initial plasma density, ultrashort pulses can be used to form the initial plasma column using optical field ionization (OFI). As was detailed in section 1.6, the fractional ionization yield from OFI only depends on the laser intensity and ionization potential of the gas. Laser-driven OFI waveguides relying on hydrodynamic shock expansion are limited by low electron temperatures comparable to the electron ponderomotive potential in the laser field at the ionization threshold of the gas [71]. At the hydrogen ionization threshold of $\sim 10^{14} \text{ W/cm}^2$, $T_e < \sim 10 \text{ eV}$ for a $\lambda = 1 \mu\text{m}$ laser, resulting in weak expansion leading to low shock walls with limited optical confinement. OFI heating can be enhanced by using a gas with higher ionization potential such as helium [72], but for a fixed pulse width, the laser energy must be increased by $> 100 \times$ compared to hydrogen.

In this chapter, we present a new approach to OFI-generated waveguides and demonstrate high intensity optical guiding up to hundreds of Rayleigh ranges. Our approach decouples the generation of the needed electron density profile from plasma heating and does not exclusively rely on hydrodynamic expansion. It leverages the idea that OFI can effectively imprint prescribed electron density profiles while minimally heating them.

2.2 Experimental Setup

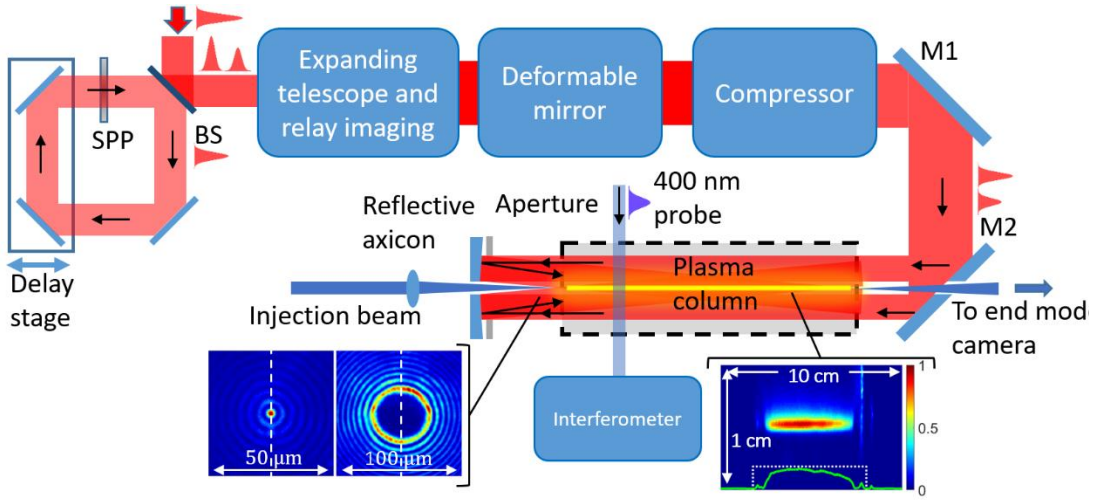


Figure 2.1. Experimental setup. The ~ 200 ps uncompressed pulse from a Ti:Sapphire amplifier is split by a non-polarizing beamsplitter (BS). The transmitted pulse is converted to a $(0, q)$ beam by a segmented q^{th} order spiral phase plate (SPP). After beam expansion and relay imaging, phase front correction using a deformable mirror, and re-compression, the beam is focused by a reflective axicon in either backfill hydrogen gas or over a supersonic gas jet. The counter-propagating guided beam is injected into the waveguide through a hole in the axicon and imaged after guiding. Also shown: frequency-doubled ($\lambda=400$ nm, 70 fs) interferometric probe pulse. Left bottom inset: images of J_0 and J_{16} beams in focus. Right bottom inset: J_0 -pulse-induced plasma fluorescence 3.5 mm above the 5 cm supersonic H₂ jet.

As detailed in Fig. 2.1, an amplified and stretched (to ~ 200 ps) linearly polarized $\lambda=800$ nm pulse from a Ti:Sapphire laser is split at an 80/20 non-polarizing beamsplitter, with 80% of the energy entering a delay ring containing a transmissive spiral phase plate (SPP) to impose either a 16π ($q = 8$) or 32π ($q = 16$) azimuthal phase screw on the beam. The recombined output beam is a $(p = 0, m = 0)$ Laguerre super-Gaussian beam pulse followed, at a delay $\tau_d = 1 - 3$ ns, by a $(p = 0, m = q)$ pulse, where p and m are radial and azimuthal mode indices. The beam is then expanded and relay imaged, phase front-corrected, recompressed to 50-100 fs FWHM, and directed to a 50.8mm diameter reflective axicon in an experimental chamber. The

$4 \times$ relay imaging conveys the plane of the spiral phase plate to the axicon. Phase front correction improves the quality of the $q = 8$ or $q = 16$ beams, employing a deformable mirror using a phase front retrieval technique developed by our group [73]. The reflective axicon forms a double-pulsed Bessel-Gauss beam, identified here using only the Bessel notation as a J_0 pulse followed $\tau_d = 1 - 3$ ns later by a J_q pulse. The J_0 pulse forms the core of the plasma waveguide by OFI; the J_q pulse forms the cladding. Here we use an H_2 backfill target or an elongated H_2 gas jet.

2.3 Extraction of the Plasma Density Profile

The plasma density and profile were measured using direct imaging of the recombination fluorescence and extracted from interferograms generated by an ultrashort $\lambda = 400$ nm, 70 fs probe pulse, imaged from the plasma through a femtosecond shearing interferometer and onto a CCD camera. Each profile was constructed using 200 shots (frames). The analysis steps are as follows. First, the 2D interferometric phase shift $\Delta\Phi(y, z)$ was extracted from each frame using standard techniques [74], where z and y are coordinates along and transverse to the optical axis. In each column (fixed z), background subtraction was performed by fitting the phase in the plasma-free part of the frame to a polynomial and interpolating over the full frame. An average phase shift $\overline{\Delta\Phi(y)}$ was then computed over 500 columns of the extracted phase and then over 200 shots. This procedure resulted in phase shift resolution of ~ 0.4 mrad, enabling extraction of electron densities as low as $4 \times 10^{16} \text{ cm}^{-3}$ (assuming a minimum signal-to-noise ratio as 20 dB). $\overline{\Delta\Phi(y)}$ was then Abel inverted to recover the plasma density.

For the symmetric plasmas generated by J_0 -pulses, the above procedure was sufficient. However, for plasmas generated by a J_q pulse, $\Delta\Phi(y,z)$ was often asymmetric. We attributed this apparent azimuthal asymmetry to asymmetry in the intensity profile of the J_q beam.

While the standard Abel inversion assumes azimuthal symmetry of the density profile to be measured, it can be used in modified form for a profile that is a separable function of radial and azimuthal coordinates (r and ϕ). Here, we use a separate measurement of the J_q intensity profile to inform the extraction algorithm. First, since the intensity of the J_q beam is near zero inside its first ring, we assume that there is no further ionization there and subtracted the phase shift $(\Delta\Phi(y,z))_{J_0}$ of the J_0 plasma from the total phase shift from the two pulses $(\Delta\Phi(y,z))_{J_0+J_q}$. We then assume that the plasma profile follows the same asymmetry as the J_q intensity profile and that its radial dependence is smooth and non-negative. A modified Abel Inversion was performed with these assumptions, and then the pre-existing J_0 plasma was added back into the electron density profile.

To account for the azimuthal asymmetry in the density profile, we added azimuthal terms to the Abel Inversion with coefficients derived from the J_q profile, and assume that the phase profile to be extracted can be expressed as $F(r,\phi) = f(r) \sum_{m=0}^n c_m \cos(m\phi + \Delta\phi_m)$. Then the plasma density profile is extracted by $\text{argmin}_F \int |\mathcal{A}[F] - \Delta\Phi(y)|^2 dy$, subjected to $F(r,\phi) \geq 0$ and $f(r < r_{min}) = 0$ ($r_{min} = 10\mu m$). Here $\mathcal{A}[F]$ is the Abel projection operator and $\Delta\Phi(y)$ is measured phase shift in the experiment. The radial profile $f(r)$ was fit by a Gaussian basis set to

improve smoothness and account for the limited resolution in interferometry, as detailed in [75]. The $\Delta\phi_m$ terms are determined by fitting the measured J_q focus profile azimuthally, as shown in Figure 2.2. In optimization we keep azimuthal terms up to $m = 6$. The optimization used the built-in MATLAB function `fmincon`, and $f(r)$ was initialized with random noise. The criterion of convergence was set as the normalized error $\epsilon = \frac{\int |\mathcal{A}[F] - \Delta\Phi(y)|^2 dy}{\int |\Delta\Phi(y)|^2 dy} < 0.001$. To further verify uniqueness of the solution, we synthesized $F(r, \phi)$ with randomly shaped ring profiles and azimuthal terms up to $m=6$, and construct Abel transform of these synthesized profiles as measured phase shift profiles. Reconstruction of these dummy phase shift profiles using this method showed correct results upon convergence when $f(r)$ was initialized by random noise. Therefore, we conclude that uniqueness is achieved.

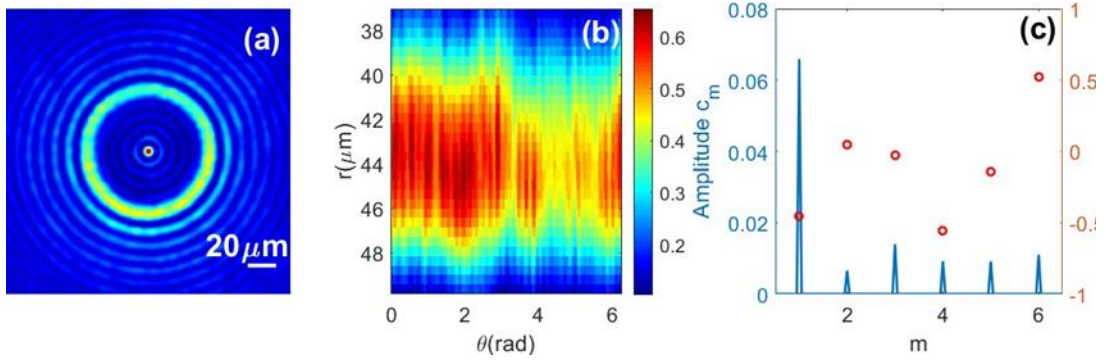


Figure 2.2. Azimuthal decomposition of a sample $J_0 + J_{16}$ focal profile. (a) Measured $J_0 + J_{16}$ focal profile. (b) Radial lineouts of main J_{16} ring vs. angle. The colorbar shows intensity normalized to the central J_0 maximum of (a). (c) Azimuthal decomposition of the radially averaged ring intensity of (b).

This extraction method produces 2D plasma density profiles, but the azimuthally averaged lineout is plotted in the rest of the chapter. The focal profile was measured along the full focal line, and the average $\Delta\phi_m$ was used in the minimization.

An example of the full 2D electron density profile is shown in Figure 2.3. The

asymmetry in Figure 2.3(a) and (c) is attributed to non-uniformity of the J_{16} beam.

The 2D plasma density profiles were used in calculation of the guided mode.

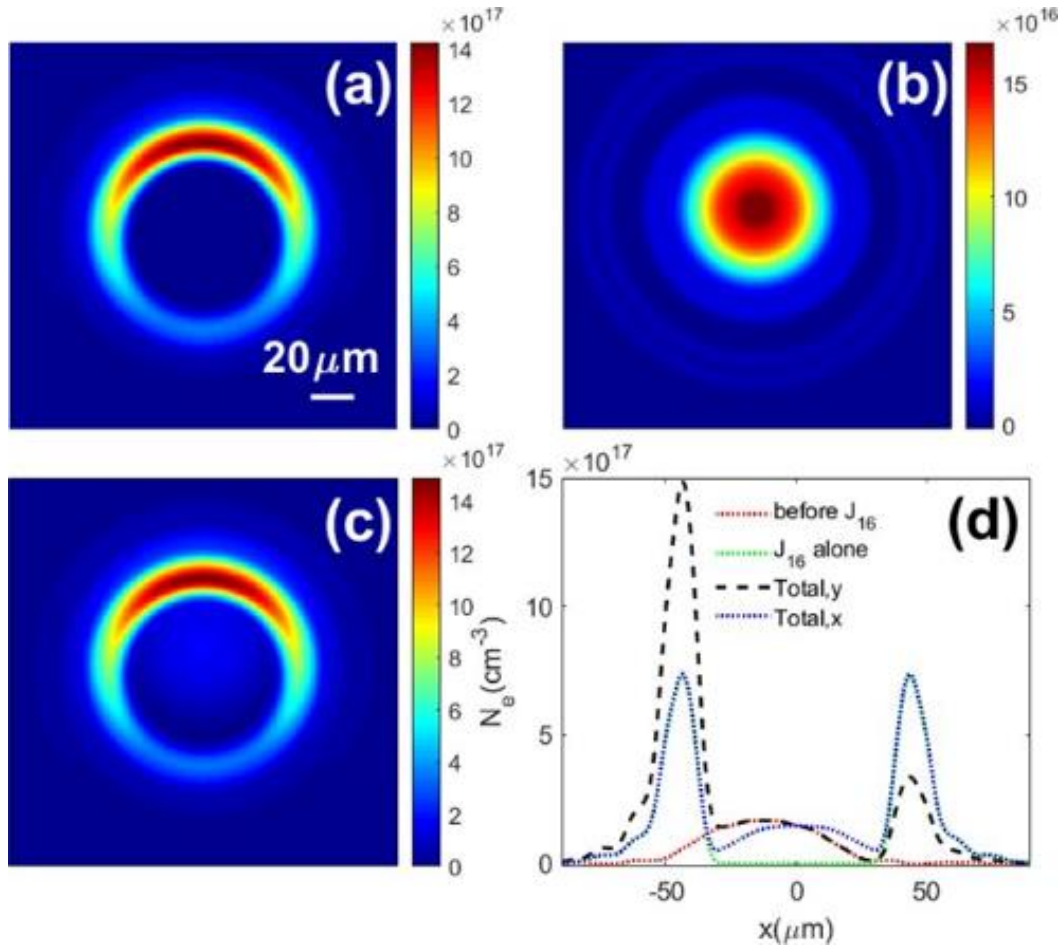


Figure 2.3 Full 2D extraction of the electron density profile shown in Fig. 2.6(f) (a) Electron density profile from the J_{16} pulse only. (b) Electron density from the J_0 pulse, right before the arrival of the J_{16} pulse. (c) Total electron density after the J_{16} pulse, (d) Lineouts of panels (a)-(c).

2.4 Waveguide Formation

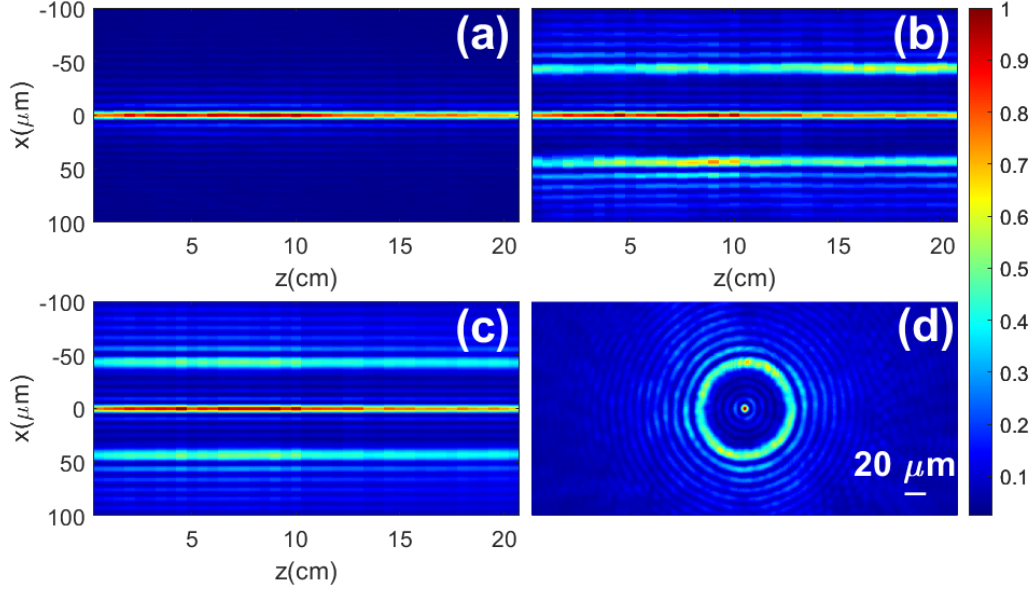


Figure 2.4. (a) Optical axis scan of J_0 focal spot profiles with planar cut shown in Fig. 2.1 inset. (b) Optical axis scan of combined J_0 and J_{16} beam, with planar cut shown in Fig. 2.1 inset. (c) Azimuthal average of (b). (d) Combined $J_0 + J_{16}$ beam profile. $\gamma = 3^\circ$.

In our experiments, the Bessel beam depth of focus, or line focus length, is $L \approx (R - a)/\tan \gamma \gtrsim 15$ or 30 cm, where $R = 2$ cm is the radius of the outer aperture near the axicon surface, a is the radius of the axicon central hole, $\gamma = 2\alpha$ is the Bessel beam ray approach angle to the optical axis, and $\alpha = 1.5^\circ$ or 3° is the axicon base angle. The inner and outer edges at a and R set up sharp transitions of Bessel beam intensity at the extremes of the focal line, preventing excessive axial taper of the plasma there. In vacuum, the central peak radius of the J_0 beam was $r_0 = 2.405/k\sin\gamma = 6$ or $3 \mu\text{m}$, where k is the laser vacuum wavenumber. The initial J_0 pulse fully ionizes the hydrogen gas target, creating a long thin plasma column which expands cylindrically outward, leaving a region of low electron density near the optical axis, creating the “core” of the plasma waveguide. The delayed J_q pulse generates a long cylindrical plasma shell [76–

78] around the expanded column: this serves as the lower refractive index “cladding”. The intense pulse to be guided is focused through a hole in the axicon and coupled into the waveguide.

An important first requirement is for the J_0 pulse to fully ionize the hydrogen gas target to create a uniform waveguide core after expansion. We found that an optical axis-averaged peak intensity of $\sim 4 \times 10^{14} \text{ W/cm}^2$ in the J_0 central maximum was sufficient to fully ionize the hydrogen gas and clamp the electron temperature along the focal line sections used in our experiments, despite the axial variation in the intensity dictated by the mapping of the beam profile to the focal line given by

$$I(z) = I_0 z \exp(-(z/w)^n) \quad (2.4)$$

with $z = r/\tan\alpha$, $r \geq a$, $n \approx 6$). Figures 2.4(a) and 2.4(b) plot cuts of J_0 alone and combined $J_0 + J_{16}$ along the white-dashed plane edges in Fig. 2.1. These composite images were obtained by scanning a $10\times$ microscope objective along z . The profiles remain relatively uniform over tens of centimeters, with z variations mostly from radial variations of the input beam. Figure 2.4(c) shows 2.4(b) azimuthally averaged, and Fig. 2.4(d) shows a transverse cut across z of the combined beams.

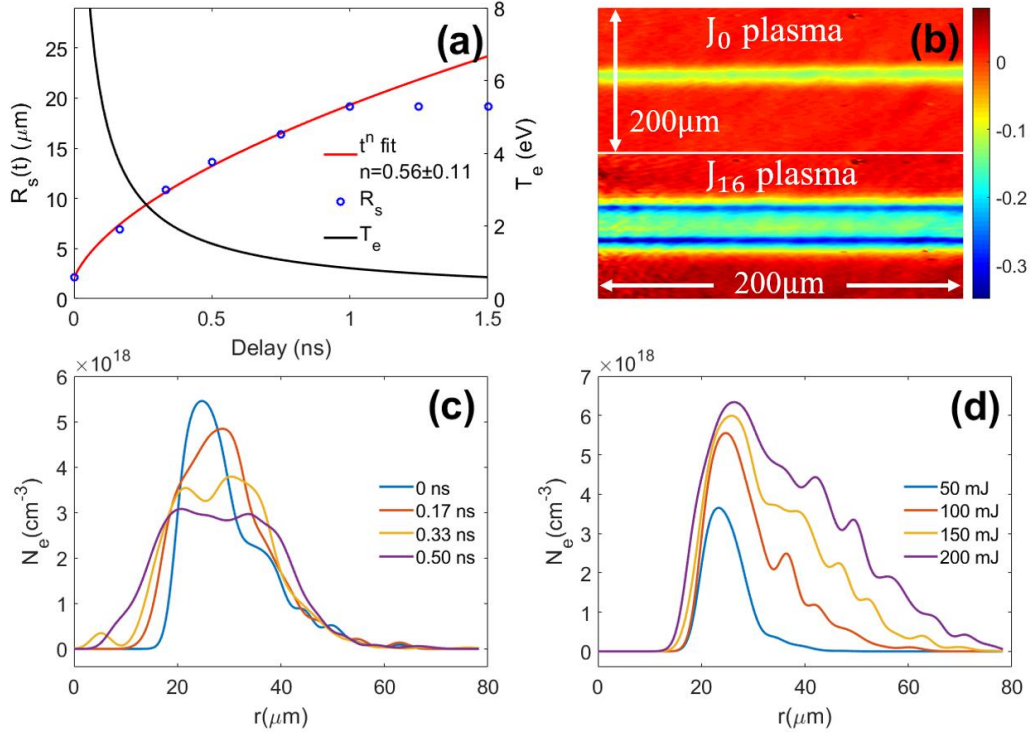


Figure 2.5. (a) Radius of blast wave expansion of J_0 plasma (best fit $n = 0.56 \pm 0.11$) and extracted electron temperature vs. time. Hydrogen backfill 50 torr. (b) Interferometric phase shift images of J_0 and J_{16} plasmas 10ps after generation. Hydrogen backfill 100 torr. (c) Time evolution of J_{16} -generated annular electron density profile. The plasma pressure gradient mainly drives filling of the central hole. Pulse energy=100 mJ. (d) Pulse energy dependence of J_{16} -generated annular electron density profiles, with saturation for >50 mJ. The modulations at larger radius are from ionization by outer rings of the J_{16} pulse.

After OFI by the J_0 pulse, the electrons thermalize over the electron-electron collision time $\tau_{ee} < 1$ ps, for $N_e \sim 10^{18} \text{ cm}^{-3}$ and initial electron temperature $T_{e0} < \sim 10 \text{ eV}$, leaving a warm plasma column surrounded by cold, neutral hydrogen gas. The plasma column rapidly cools as it expands outward as a cylindrical blast expansion $R_s(t) = \xi_0(\varepsilon_0/\rho)^{1/4}t^{1/2}$, where R_s is the radial position of the outer plasma boundary, ε_0 ($\propto T_{e0}$) is the laser energy deposition per unit length, ρ the initial mass density of the hydrogen gas, and ξ_0 is a dimensionless prefactor of order unity [79]. Typical plots

of $R_s(t)$ and associated density profiles are shown in Fig. 2.5(a) and Fig. 2.6(a) for peak J_0 intensity 10^{15} W/cm², and backfill hydrogen pressure of 50 Torr. The electron temperature T_e is determined from the ion acoustic speed as $dR_s/dt \approx c_s = (\gamma_c T_e / m_i)^{1/2}$, where γ_c is the specific heat ratio and m_i is the proton mass. It is seen that after a few nanoseconds, the expansion stagnates as the plasma cools to < 1 eV, with the central plasma density reduced by $\sim 10 \times$ and the column expanded by $\sim 3 - 5 \times$. Two-color interferometry and integration of the total phase shift from the plasma column indicate that there is no recombination on this time scale. The weak $(T_{e0})^{1/4}$ dependence of the plasma expansion ensures that moderate axial variations in Bessel beam intensity negligibly affects the local axial uniformity of the waveguide plasma, as seen in Fig 2.5(b).

After the J_0 pulse has formed the core of the plasma waveguide, the delayed J_q pulse forms the cladding. The index q of the high order beam and its delay are chosen so that it propagates through the expanded J_0 plasma column, forming a high intensity annulus at its edge, ionizing the neutrals there. The annular plasma formed by a J_{16} pulse alone is shown in Fig. 2.5(c), as function of delay after generation at $t = 0$, and in Fig. 2.5(d) at $t = 0$ as a function of J_{16} pulse energy. It is seen that the annular plasma cladding decays over ~ 0.5 ns, mainly by filling in the central hole, and that the cladding height saturates at > 50 mJ.

The effect of the J_0 –generated plasma on the propagation of the J_q pulse was calculated using a 3D UPPE simulation [80] implementation called YAPPE (‘Yet Another Pulse Propagation Effort’). UPPE (‘Unidirectional Pulse Propagation Equation’) is a system of ordinary differential equations (ODEs) of the form

$$\frac{\partial}{\partial z} A_{k_x, k_y}(\omega, z) = iQ_{k_x, k_y}(\omega) 2\pi P_{k_x, k_y}(\omega, z) e^{-i\left(k_z - \frac{\omega}{v_g}\right)z}. \quad (2.5)$$

$A = A_{k_x, k_y}(\omega, z)$ is an auxiliary field related to the Fourier transform of the optical field by $E = Ae^{ik_z z}$. The spectrum of transverse spatial frequencies (k_x, k_y) indexes a system of ordinary differential equations, which is solved using a GPU implementation of MATLAB's ODE45 function. $P_{k_x, k_y}(\omega, z)$ is the nonlinear polarization of the medium including (for our simulations in this paper): Kerr self-focusing, ionization losses, and a non-dispersive plasma response. The ionization rate is computed using an ADK module with neutral density saturation [21,81]. The perturbed index of refraction due to the H₂ neutral density gradient is also included as a part of the nonlinear polarization. The nonlinear refractive index (n_2) of hydrogen [82] is scaled with the spatiotemporally variable H₂ neutral density such that it goes to zero when all neutrals are ionized.

The higher order Bessel (J_q) beams were initialized as LG_{0q} modes with an applied axicon phase. The axicon phase was defined in YAPPE as $\varphi = -k_0 r \tan(\alpha)$ for central wavenumber k_0 , distance from central axis r , and axicon base angle α . Figure 2.6(b) shows the simulated J_{16} pulse-generated electron density profiles with and without the J_0 -generated core plasma present. The effect of the preformed core plasma on the J_{16} -pulse propagation and induced OFI profile is seen to be negligible. In general, unless the core plasma density approaches the effective critical density $N_{cr,eff} = N_{cr} \sin^2 \gamma \sim 0.5 - 2 \times 10^{19} \text{ cm}^{-3}$ experienced by oblique Bessel beam rays, it has a negligible effect on the J_q -induced ionization profile.

2.5 Guided Mode Analysis

The two Bessel-pulse method of plasma waveguide generation enables very wide tuning of the plasma waveguide transverse profile and its guided modes through control of the hydrogen density, the order q of the second pulse, the Bessel beam axis approach angle γ ($= 2\alpha$), the energy and pulsewidth of the two pulses, and the time delay between them. The plots in Fig. 2.6(c)-(f) illustrate this flexibility.

In general, the desired guided optical mode dictates the required waveguide profile. As discussed in detail in [83], plasma waveguides with a finite thickness cladding have modes that are quasi-bound or leaky. A sufficiently well-bound leaky mode has an exponential decay length that is long compared to other scale lengths of interest, such as the dephasing length or pulse decay due to energy depletion in laser wakefield acceleration. There are a variety of methods for calculating leaky modes, and we used two different methods for azimuthally symmetric 1d profiles and asymmetric 2d profiles.

Leaky modes were found for azimuthally averaged density profiles by solving the Helmholtz equation for the field $E(r, z) = \mathcal{E}(r)e^{i\beta z}$ in cylindrical coordinates,

$$d^2\mathcal{E}/ds^2 + s^{-1} d\mathcal{E}/ds + (n^2 - \beta^2/k_0^2)\mathcal{E} = 0 \quad (2.6)$$

where k_0 is the vacuum wave number, β is the waveguide propagation wavenumber, $s = k_0 r$, and $n(r)$ is the refractive index profile derived from the azimuthally averaged plasma profile. To map out the quasi-bound modes for a given index profile $n(r)$, we repeatedly solve Eq. (2.5), scanning $\beta' = \beta/k_0$ and plotting $\eta(\beta') = (|\mathcal{E}_{vacuum}|^2 A)^{-1} \int_A |\mathcal{E}|^2 dA$, where the integral is over the waveguide cross section. The main peaks in $\eta(\beta')$ identify the quasi-bound modes, with the full width at half

maximum $\Delta\beta$ of these peaks giving the axial (z) $1/e$ attenuation length of intensity $L_{1/e} = \Delta\beta^{-1}$.

Guided modes for the full 2D plasma profiles were calculated using a Beam Propagation Method (BPM) simulation [84]. A test mode overfilling the plasma density profile was input into the waveguide and propagated using a split-step Fourier method over more than 100 Rayleigh lengths until all the high-order modes have decayed away. The fundamental guided mode is then propagated in the same manner and the peak intensity is sampled and fit to an exponential function to determine $L_{1/e}$.

A good design approximation that requires much less analysis than those two methods recognizes the step-like profiles of our two-pulse OFI waveguides: a relatively flat core and sharply rising cladding, for which the fundamental channel mode radius is given by [85]

$$w_{ch} \approx a(0.6484 + 1.619V^{-3/2} + 2.879V^{-6} + \dots) \quad (2.7)$$

where a is the core radius and the step index fiber parameter is $V = ka(n_{core}^2 - n_{cladding}^2)^{1/2} = a(4\pi r_e \Delta N_e)^{1/2}$, where $\Delta N_e = N_e^{cladding} - N_e^{core}$, r_e is the classical electron radius, and k is the laser vacuum wavenumber. Larger guided modes can be supported by smaller core-cladding differences ΔN_e , which require lower hydrogen density N_0 (since under OFI, ΔN_e is roughly proportional to N_0).

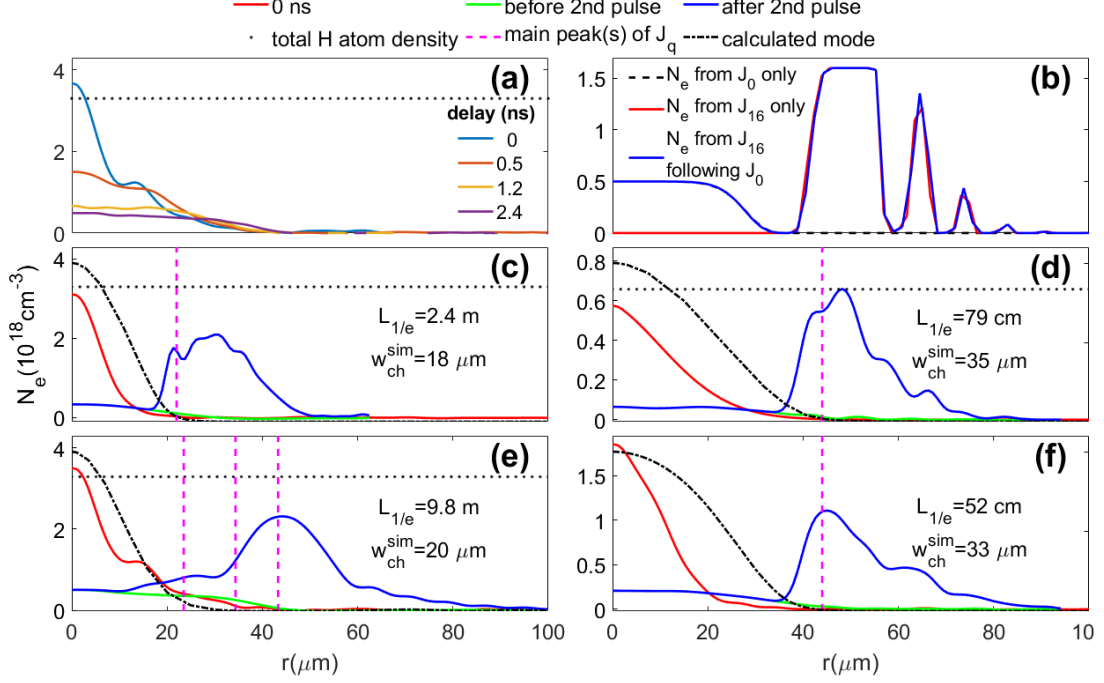


Figure 2.6. Transverse electron density profiles in backfill and gas jet hydrogen plasmas. (a) J_0 -generated ‘core’ electron density profiles vs. time delay (b) J_{16} pulse propagation simulations of plasma profile generation with and without J_0 -generated core plasma present. (c) – (f): Electron density profiles measured ~10 ps after the J_0 pulse (red curves), and 10 ps before and after the J_q pulse (green curves and blue curves). (c) 100 mJ (J_0) +235 mJ (J_{16}), $\alpha = 3^\circ$, H2 backfill pressure $P=50$ torr. Waveguide length $L = 15$ cm; (d) 78 mJ (J_0)+250 mJ (J_{16}), $\alpha = 1.5^\circ$, H2 backfill pressure $P=10$ torr. Waveguide length $L = 30$ cm; (e) 100 mJ (J_0)+235 mJ (J_8), $\alpha = 1.5^\circ$, H2 backfill pressure $P=50$ torr. Waveguide length $L = 30$ cm; (f) H2 gas jet, $P \approx 30$ torr: 87 mJ (J_0)+235 mJ (J_{16}), $\alpha = 1.5^\circ$, $q=16$. Waveguide length $L = 5$ cm.

This mode scaling is illustrated in Fig. 2.6, where the H₂ density in Fig. 2.6(c) is chosen to be $5 \times$ higher than in Fig. 2.6(d). Both experiments use a J_{16} pulse to generate the waveguide cladding, the approach angle decreasing from $\gamma = 6^\circ$ in Fig. 2.6(c) to $\gamma = 3^\circ$ in Fig. 2.6(d). This has the effect of widening the initial J_0 -generated fully ionized core plasma and moving the cladding peak from a radial position of $\sim 30 \mu\text{m}$ out to $\sim 50 \mu\text{m}$. In both cases, the core plasma remains unaffected by the second pulse, as seen by the merging of the blue and green curves toward the waveguide axis. In Fig.

2.6(c), $\Delta N_e \sim 1.3 \times 10^{18} \text{ cm}^{-3}$, $a \sim 20 \mu\text{m}$, and $V \sim 4.3$, predicting $w_{ch} \sim 17 \mu\text{m}$, while in Fig. 2.6(d), $\Delta N_e \sim 0.5 \times 10^{18} \text{ cm}^{-3}$, $a \sim 40 \mu\text{m}$, and $V \sim 5.4$ give $w_{ch} \sim 32 \mu\text{m}$.

We also plot the fundamental modes and their $1/e$ intensity attenuation lengths $L_{1/e}$ computed using our leaky mode solver, showing good agreement with the step index values. As an example of the control available for sculpting transverse profiles, Fig. 2.6(e) shows use of a J_8 pulse at approach angle $\gamma = 3^\circ$. Here, the peaks of the first and second J_8 rings overlap with $\sim 50\%$ ionized hydrogen in the expanded plasma from the J_0 pulse, generating an inner cladding bump, while the third ring overlaps neutral hydrogen gas on the plasma periphery, generating a higher outer cladding bump. Considering this as a step index profile to the first cladding bump (with $\Delta N_e \sim 0.5 \times 10^{18} \text{ cm}^{-3}$, $a \sim 15 \mu\text{m}$, and $V \sim 2.0$), gives $w_{ch} \sim 19 \mu\text{m}$, in good agreement with the leaky mode solver.

We have presented results thus far for waveguides generated in H_2 backfill gas. For LWFA, however, distortion-free injection and guiding of high intensity pulses demands that the waveguide entrance and exit be free of neutral gas. For this purpose, we have developed a supersonic gas jet producing uniform low-pressure flows. The inset to Fig. 2.1 shows uniform fluorescence from J_0 pulse induced OFI of a 5 cm long H_2 gas sheet, 3.5 mm above the nozzle orifice where the pressure is ~ 30 Torr. The 15 cm long Bessel beam focus overfills the gas sheet. For jets longer than ~ 5 cm, the gas nozzle interferes with J_q formation in the focus; we have recently developed an optical fix for this limitation [86]. Figure 2.6(f) shows results from 2-pulse OFI waveguide generation in the gas jet, where a step-like guide is clearly formed. Applying the step index analysis

(using $\Delta N_e \sim 0.75 \times 10^{18} \text{ cm}^{-3}$, $a \sim 40 \mu\text{m}$, and $V \sim 2.0$) gives $w_{ch} \sim 30 \mu\text{m}$, in reasonable agreement with the leaky mode computation.

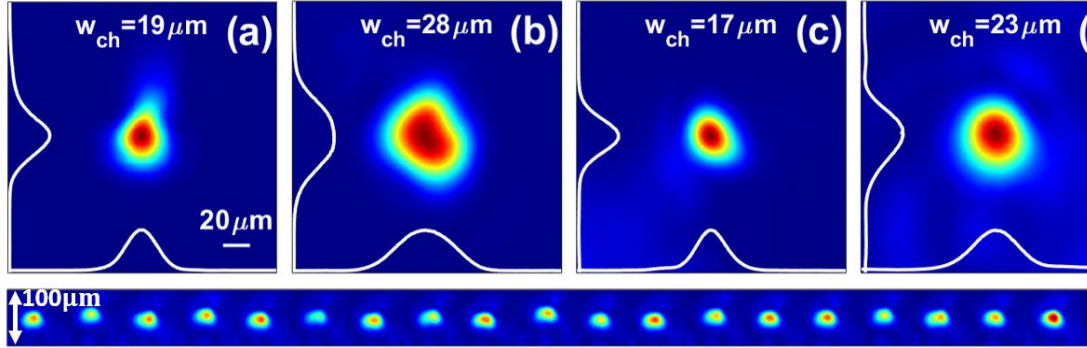


Figure 2.7. (a)(b)(c) Low intensity exit modes from $L = 15$, 30 , and 30 cm long plasma waveguides (number of Rayleigh ranges $L/z_0 \sim 110$, 100 , and 260 (for $\lambda = 800 \text{ nm}$)) under conditions of Fig. 2.6(c)(d)(e), for injection of $\sim 100 \mu\text{J}$, $\lambda = 400 \text{ nm}$ pulse. (d) Gas jet experiment: high intensity exit mode of a 5 cm long waveguide generated in H_2 gas jet. Injected pulse 40 mJ , $\lambda = 800 \text{ nm}$, 50 fs FWHM, with $\sim 50\%$ coupling efficiency, giving guided intensity $\sim 6 \times 10^{16} \text{ W/cm}^2$ over $L/z_0 \sim 25$ Rayleigh ranges (e) Guided modes from 20 successive gas jet shots for conditions of (d). Fluctuations are from beam pointing fluctuations of waveguide formation and injected beams.

Guided modes for the $L = 5 \text{ cm} - 30 \text{ cm}$ waveguides of Fig. 2.6 are shown in Fig. 2.7, where the modes of Fig. 2.7(a)-(d) correspond to the conditions of Fig. 2.6(c)-(f) (and (2-7(d)(e) and 4(f) are for the jet experiment). For these spot sizes and guide lengths, the number of Rayleigh lengths of guiding (at $\lambda = 800 \text{ nm}$) ranges over $L/z_0 \sim 25$ to ~ 260 . In all cases, backfill and gas jet, the guided pulse was injected into the waveguide within $\sim 10 \text{ ps}$ after the cladding-generating J_q pulse. Even though the core expands over $\sim 2 \text{ ns}$ and the cladding collapses in $\sim 0.5 \text{ ns}$, as seen in Figs. 2.4 and 2.5, this does not limit the length of fixed-parameter waveguides produced using our method: any guided pulse can propagate a couple of nanoseconds behind the J_0 pulse and immediately behind the traveling wave cladding formation by the J_q pulse. As can be seen in Fig. 2.7, the guided modes are slightly asymmetric owing to imperfectly

corrected azimuthal intensity variations in the J_q beam. This leads to slight azimuthal guide variations, as we have determined using the quasi-2D Abel inversion algorithm employing the measured J_q beam asymmetry. The guided mode radii w_{ch} in the figure are from 2D Gaussian fits to the measured mode and agree reasonably well with the computed modes in Fig. 2.6(c)-(e) (the guide in 2.6(f) may be tapered at the end), where the electron density profiles were azimuthally averaged from the extracted 2D profiles.

The guided mode in the hydrogen jet plasma waveguide (Fig. 2.7(d)) has peak intensity $\sim 6 \times 10^{16} \text{ W/cm}^2$, based on $\sim 50\%$ coupling efficiency, where the injected laser energy was limited to 40 mJ by our system's available laser energy. There was no evidence of additional ionization by the guided pulse, consistent with our density profile measurements showing fully ionized waveguide cores. Exit modes from a sequence of 20 successive gas jet shots is shown in Fig. 2.7(e). Given the pointing sensitivity of the waveguide forming and injection pulses, the shot-to-shot reproducibility is reasonable. Future versions of this experiment will require active stabilization of all beams.

Based on the results presented here, we can estimate potential electron acceleration driven by a more powerful laser, which would enable longer waveguides and higher injection energies. As an example, the waveguide in Fig. 2.6(d) has a core electron density of $N_{e0} \sim 5 \times 10^{16} \text{ cm}^{-3}$ and $w_{ch} = 35 \mu\text{m}$. This corresponds to $L_{deph} \approx 1.3 \text{ m}$ (Eq. 1.10), which exceeds $L_{1/e} \sim 0.8 \text{ m}$ for this guide. For a temporally resonant pulse with $a_0 = 1.15$ (for which $L_d \approx L_{pd}$), we can apply Eq. 1.9 for E_z and find the energy gain $\Delta W \approx 7 \text{ GeV}$. This example requires laser energy 6 J and pulse length 100 fs for a peak power of 55 TW. Improving the mode confinement so $L_{1/e} > L_d$ would give

$\Delta W \sim 11 \text{ GeV}$. Increasing the guided spot size to $w_{ch} \approx \lambda_p/2 \sim 75 \mu m$ for the same N_{e0} , and extending the corresponding waveguide so that $L_{1/e} > L_d = 2.1 \text{ m}$, would yield $\Delta W \sim 34 \text{ GeV}$ for laser energy 21 J and peak power 190 TW. For this case, the laser energy cost for the waveguide would be $< \sim 5 \text{ J}$.

In conclusion, the two Bessel method is a new, highly tunable technique for generating long, low loss plasma waveguides for a wide range of plasma densities and mode sizes. We have demonstrated guiding in waveguides up to 30 cm and over 250 Rayleigh lengths, with a guided intensity limited only by our available laser energy.

Chapter 3: Self-Waveguiding

3.1 Introduction

In this chapter, we investigate another, quite different approach: the “self-waveguiding” of relativistic laser pulses in prepared neutral gas refractive index structures. The physics of self-waveguiding is quite distinct from the well-known process of relativistic self-guiding described in section 1.5, where relativistic self-focusing leads to electron cavitation by the laser ponderomotive force [14,87]. In self-waveguiding, the front of the pulse generates a real plasma waveguide that confines the rest of the pulse—the front of the pulse does the optical field ionization ‘work’ of generating the waveguide, and then the rest of the pulse is channeled by it. Relativistically intense pulses are not required, but pulses must exceed the optical field ionization threshold of the target gas. The process requires an initially prepared elongated refractive index structure whose necessary feature is a gas density minimum on axis. This requires impulsive heating of the gas along an extended path; we show that plasma generation by optical field ionization of hydrogen gas by a Bessel beam pulse supplies sufficient heating. An initial plasma is necessary only to the extent that it heats the gas to generate the axial neutral density depression. We note that prior work by Morozov *et al.* [88] had shown that the leading edge of a pulse injected into an incompletely ionized plasma waveguide can ionize the peripheral neutrals and lead to guiding over a few mm, and that laser pulses injected into capillary discharge waveguides can modify them to improve guiding [89,90].

Here, we present detailed experiments and simulations exploring self-waveguiding of weakly relativistic pulses over 10 cm (100 Rayleigh ranges) and measuring in detail

the guided mode structure and time resolved refractive index profiles of plasma and neutral hydrogen. We show that during self-waveguiding, guided mode beating can be significant. A self-waveguided pulse can also be viewed as a vehicle for generating the plasma waveguide that is left in its wake, suitable for guiding an even more intense follow-on pulse to drive applications such as LWFA. To demonstrate this principle, we demonstrate guiding of a probe pulse that closely follows the self-waveguided pulse.

3.2 Experimental Setup

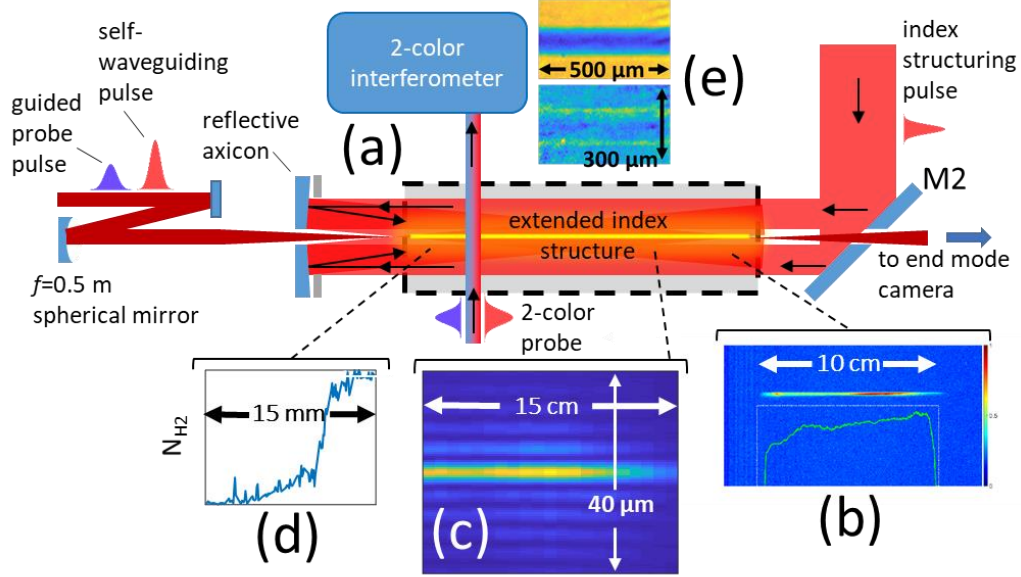


Figure 3.1. (a) Experimental setup, showing index structuring Bessel beam pulse (J_0 pulse) (100 mJ, 75 fs linearly polarized, $\lambda=800$ nm) focused by a reflective axicon with base angle $\alpha = 3^\circ$, over a 10 cm long supersonic hydrogen gas jet. After a delay Δt_{inj} , a 40 fs, 10-90 mJ, linearly polarized pulse (self waveguiding pulse) is focused through a hole in the axicon at $f/40$ and injected into the entrance of the index structure. The exit of the index structure is imaged by an end mode CCD camera. A follow-on pulse (guided probe pulse) is guided by the plasma waveguide generated by the self-waveguiding pulse. (b) H₂ gas density profile and fluorescence from plasma recombination after optical field ionization by the J_0 pulse. (c) Magnified axial scan of the J_0 beam profile. (d) ~2mm density falloff at the entrance of the jet. (e) 2-colour ($\lambda=400$ nm and 800 nm) interferometric probe pulse is directed transversely through the index structure and into a folded wavefront interferometer, with extracted phase maps from the electron density (top) and neutral (bottom) profiles shown.

Figure 3.1 shows the experimental setup. A 75 fs, 100 mJ, 10 Hz Ti:Sapphire pulse (index structuring pulse) is focused by a reflective axicon with a $d_h = 1.27$ cm diameter central hole and base angle $\alpha = 3^\circ$ (Fig. 3.1(a)) over a 10-cm long pulsed supersonic H₂ gas jet (Fig. 3.1(b)), generating a refractive index structure composed of plasma and neutral hydrogen. The length of the resulting Bessel beam (J_0 pulse) focus is $L_B \approx (D - a_h)/2\tan(2\alpha) \approx 16$ cm, where $D = 4.5$ cm is the beam diameter,

overfilling the jet 0.8 mm above its 10 cm long slot orifice, where the local hydrogen molecular density was adjusted in the range $N_{H_2} = 10^{18} - 10^{19} \text{ cm}^{-3}$.

A central 2D slice of the axially scanned J_0 pulse focal image is shown in Fig. 3.1(c). After a delay Δt_{inj} , a 40 fs, 10-90 mJ, linearly polarized pulse (self waveguiding pulse) is focused through a hole in the axicon at $f/40$ and injected into the entrance of the refractive index structure. The exit of the index structure is imaged by an end mode CCD camera. For experiments in a pump-probe configuration, a slightly delayed probe pulse is injected into the guide generated by the self-waveguiding pulse. The gas sheet is generated by a 10 cm long supersonic slit nozzle (Mach ~ 4), with hydrogen gas fed into the nozzle by fast pulsed valves. The valve open time was 3 ms to ensure equilibrium gas flow. The gas density fall-off scale length at the ends of the gas sheet is ~ 2 mm (Fig. 3.1(d)) A 2-colour interferometric probe pulse ($\lambda=400$ nm and 800 nm) passes transversely through the index structure and into a folded wavefront interferometer, enabling separate extraction [91] of the neutral hydrogen and electron density contributions to the refractive index profile, with sample phase maps shown in Fig. 3.1(e). Knowledge of the individual contributions to the composite index profile is crucial for engineering the self-waveguiding process.

3.3 Results and Discussion

To begin, we assess the ability of an OFI-heated hydrogen plasma to confine and guide a low intensity end-injected laser pulse by examining the plasma's transverse refractive index profile. The 2-colour interferometric probe pulse was passed transversely through the J_0 pulse-induced index structure 1 cm past its entrance (the end closest to the axicon). Figure 3.2(a) shows time resolved electron density and

neutral hydrogen density profiles from Abel-inversion of extracted interferometric phase images, where the error bars, here and in subsequent plots, are from phase error propagation in Abel inversion. Just after arrival of the J_0 Bessel beam pulse at $t = 0$, the electron density profile is peaked on axis at $\sim 3.5 \times 10^{18} \text{ cm}^{-3}$, after which the peak drops to $\sim 5 \times 10^{17} \text{ cm}^{-3}$ over 4 ns as the profile radially expands. For our 100 mJ J_0 pulse and 50 torr hydrogen target, the electron density profile never develops an on-axis minimum and shock wall cladding characteristic of a hydrodynamic plasma waveguide [52,58,92–94], and thus it would be unable to optically guide a laser pulse. The neutral hydrogen density profile at $t = 0$ (obtained from $N_0 - N_e(r, t = 0)$) is depleted on axis and at later times develops a steepened annular enhancement as peripheral neutral hydrogen is “snowplowed” and shock-compressed by the cylindrically expanding electron density column. The peak neutral density in this annular shock wall is $> \sim 2N_0$ (and $> \sim 3N_0$ for 100 torr hydrogen) and $< \sim 20N_{e0}$. The electron density column expands and cools as a cylindrical blast wave expansion $\langle r_{\text{plasma}} \rangle = \xi_0(\varepsilon_0/\rho)^{1/4}t^{1/2}$, where $\langle r_{\text{plasma}} \rangle$ is the density-weighted average plasma radius, ε_0 ($\propto T_{e0}$) is the J_0 pulse energy deposition per unit length, ρ is the initial mass density of the hydrogen gas, and ξ_0 is a dimensionless prefactor of order unity. Figure 3.2(b) plots the time evolution of $\langle r_{\text{plasma}} \rangle$ and the peak of the neutral annulus, with a fit to $\langle r_{\text{plasma}} \rangle \propto t^n$ giving $n = 0.5 \pm 0.4$. The plasma temperature, also plotted in Fig. 3.2(b), is calculated from $(d\langle r_{\text{plasma}} \rangle/dt) \approx (\gamma_c T_e/m_i)^{1/2}$, where γ_c is the specific heat ratio and m_i is the proton mass. After a few nanoseconds, the plasma expansion stagnates as it cools to $> \sim 1 \text{ eV}$, with its central density reduced by $\sim 10 \times$ and the column expanded by $\sim 5 \times$. The neutral annulus speed is $v_N =$

7.7×10^5 cm/s (Mach ~ 6 in hydrogen [95])), giving a neutral hydrogen shock temperature $T_H \sim 0.4$ eV, consistent with the approach to electron-ion local temperature equilibration. The neutral annulus is clearly a neutral hydrogen shock wave that separates from the stagnating plasma and propagates outward at its local sound speed. We simulate the generation of a neutral hydrogen shock initiated by a central 10 eV hydrogen plasma, using a 1D Lagrangian hydrocode described in [96], with shock position vs time overlaid on Fig. 3.2(b). The computed shock velocity, $\sim 7 \times 10^5$ cm/s, is consistent with the measurements.

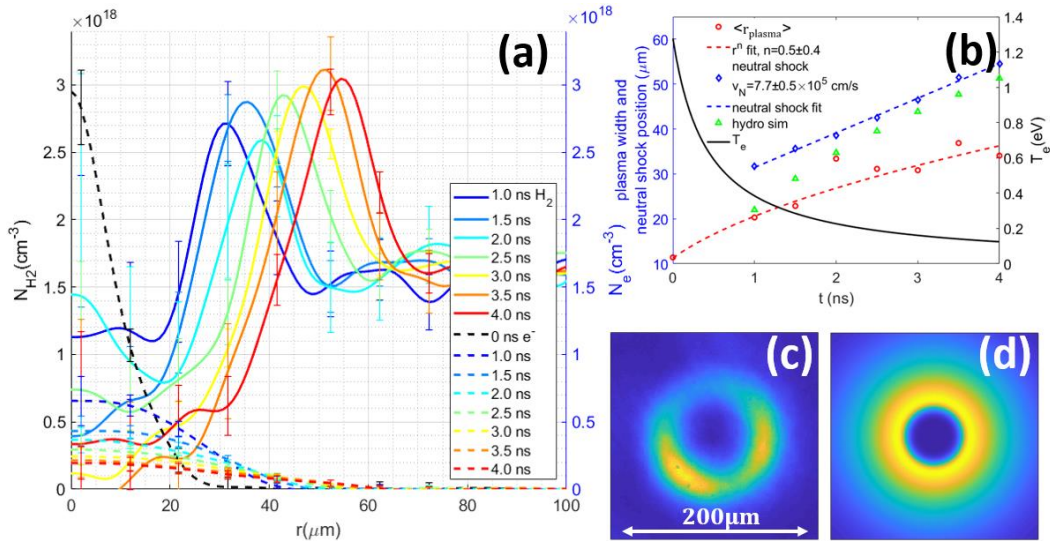


Figure 3.2. (a) Plasma and neutral gas density profile evolution measured by 2-color transverse interferometry. (b) J_0 pulse-induced RMS plasma column radius $\langle r_{\text{plasma}} \rangle$ and neutral shock position vs time. The plasma temperature is computed from the plasma expansion (see text). The neutral shock speed $v_N \sim 7.7 \times 10^5$ cm/s, calculated from the plotted points, indicates a gas temperature ~ 0.4 eV. Green triangles: neutral shock position vs. time from hydrodynamic simulation. (c) Ring mode image of $\lambda=400$ nm probe pulse trapped and guided in neutral shock cylindrical annulus. Hydrogen backfill, 50 torr. (d) Beam propagation method (BPM) simulation [84] using conditions of (c). Error bars, here and in subsequent plots, are from phase error propagation in Abel inversion.

Another way to image the annular neutral shock wall is to guide an axial probe pulse in it; we have discussed this previously in the context of cylindrical single-cycle acoustic fronts shed by femtosecond filaments in air [97]. Here, at $\Delta t_{inj} = 2.5$ ns delay after the J_0 pulse, we injected a $\lambda=400$ nm, <100 μ J probe pulse at the entrance of the 10 cm long structure and imaged the guided mode at the exit, with the result being the ring mode shown in Fig. 3.2(c). This is in good agreement with beam propagation method [84] simulations, shown in Fig. 3.2(d), for $\lambda=400$ nm Gaussian beam injection into the measured 2.5 ns delay profile of Fig. 3.2(a): the ring mode is laser light trapped by the enhanced density of the cylindrical shock wave.

At first glance, the refractive index structure described so far—a long, low density plasma cylinder surrounded by a cylindrical annulus of enhanced neutral density—appears unsuitable as a guiding structure for high intensity laser pulses. However, the ability of the cylindrical shock wave to guide a low intensity pulse (Fig. 3.2(c)) ensures that the transverse wings of the leading edge of a high intensity pulse will be guided. Later time slices of the pulse envelope might then ionize the neutral annulus, generating—on the fly—a suitable plasma waveguide cladding for the remainder of the pulse envelope.

Results of an experiment to test this scenario are shown in Fig. 3.3. Here, a high-power ($P < 5$ TW, linearly polarized, $\lambda = 800$ nm, $\tau_{FWHM} = 40$ fs) laser pulse was focused by a 0.5 m focal length spherical mirror (at $f/40$) to a $w_0 = 25$ μ m spot (Rayleigh range $z_0 = \pi w_0^2 / \lambda \sim 2.5$ mm) at the entrance of the index structure generated by the J_0 pulse. Here, the uncorrected focal spot had $\sim 30\%$ energy within a radius w_0 . The maximum peak intensity on axis is $I_0 = 2 \times 10^{17}$ W/cm² and falls to

$\sim 10^{14} \text{ W/cm}^2$ at a radial distance $r \sim 50 \text{ }\mu\text{m}$, sufficient for the far transverse wings to ionize any of the annular neutral gas profiles plotted in Fig. 3.2(a). The electron density profiles in Figure 3.3(a) show the effect of increasing the energy of a pulse injected at a delay $\Delta t_{inj} = 2.5 \text{ ns}$ (with respect to the J_0 pulse) into the refractive index structure shown in Fig. 3.2(a). Here, the interferometric probe pulse sampled the plasma 2 cm downstream ($\sim 8z_0$) from the structure entrance. It is seen that as pulse energy increases, a plasma waveguide structure forms and converges to a near-final form, consistent with the initial neutral density annulus, except for slight increases in cladding height and width. For fixed injected pulse energy (88 mJ) and $\Delta t_{inject} = 2.5 \text{ ns}$, Fig. 3.3(b) shows the increasing plasma density of the cladding as gas pressure is increased. After intense pulse injection, the annular cylinder of new plasma is driven inward and outward by the pressure gradients on the inside and outside of the annulus, distorting the guiding structure, as seen in Fig. 3.3(c). The inward pressure wave even causes plasma enhancement (compression) on axis, as seen in prior work [98]. The effect of end-injection at various delays after the J_0 pulse is shown in Fig. 3.3(d)— at increasing Δt_{inj} , the annular zone of additional ionization moves radially outward, following the location of the neutral shock.

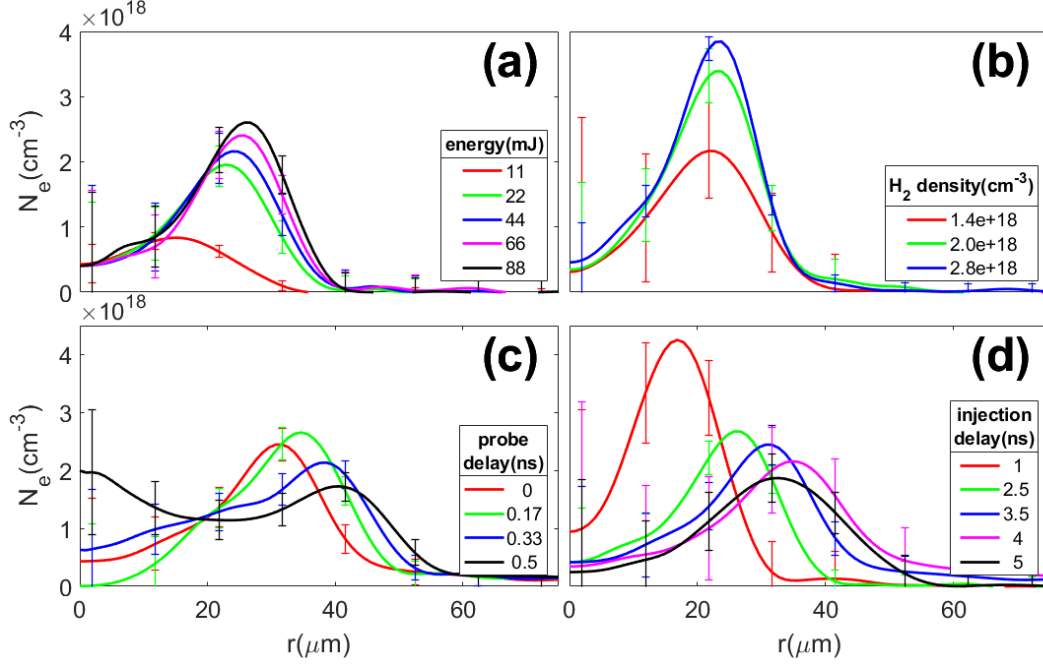


Figure 3.3. (a) Electron density profiles generated at $z = 2$ cm vs injected pulse energy. $\Delta t_{inj} = 2.5$ ns, $N_{H_2} = 1.6 \times 10^{18} \text{ cm}^{-3}$. (b) $\Delta t_{inj} = 2.5$ ns, injection beam energy = 88 mJ. (c) $\Delta t_{inject} = 3.5$ ns, $N_{H_2} = 1.6 \times 10^{18} \text{ cm}^{-3}$, injection beam energy = 88 mJ. (d) $N_{H_2} = 1.6 \times 10^{18} \text{ cm}^{-3}$, injection beam energy = 88 mJ. All panels: Probe delay (with respect to injection pulse) = 5 ps.

While Fig. 3.3 clearly shows that the injected pulse transforms the initial refractive index profile into a suitable plasma waveguide structure at an axial location $8z_0$ past the entrance, the pulse has evidently undergone guided propagation to that point and beyond. To obtain physical insight into the propagation and ionization dynamics, we performed simulations of end-injection of intense pulses into a range of refractive index profiles. We used YAPPE, described in more detail in section 2.4 where it was applied to higher order Bessel beam propagation through the plasma column. Here YAPPE is applied to an intense $\lambda=800$ nm, $\tau = 40$ fs FWHM Gaussian pulse of adjustable waist w_0 ($1/e^2$ intensity radius) placed at the structure entrance at $z = 0$. For the weakly relativistic pulses considered here, the effects of the laser ponderomotive force on the

refractive index structure are minimal, so an UPPE approach is physically justified and computationally efficient.

Figures 3.4 and 3.5 show simulations for peak intensities $I_0 = 3 \times 10^{17} \text{ W/cm}^2$ and $1.5 \times 10^{17} \text{ W/cm}^2$, respectively, injected at the structure entrance at $z = 0$. These intensities bracket our estimated peak injected intensity of $2.2 \times 10^{17} \text{ W/cm}^2$ and, as we will see, illustrate the intensity sensitivity of the self-waveguiding process. We note that simulations of even higher intensities would require particle-in-cell simulations: as $a_0 \rightarrow 1$, the laser ponderomotive force begins to modify the electron density profiles. Our goal for this chapter's simulations is to study the nonlinear optics of self-waveguiding and model our experiments for $a_0 < \sim 0.3$, for which an UPPE-based model is quite adequate.

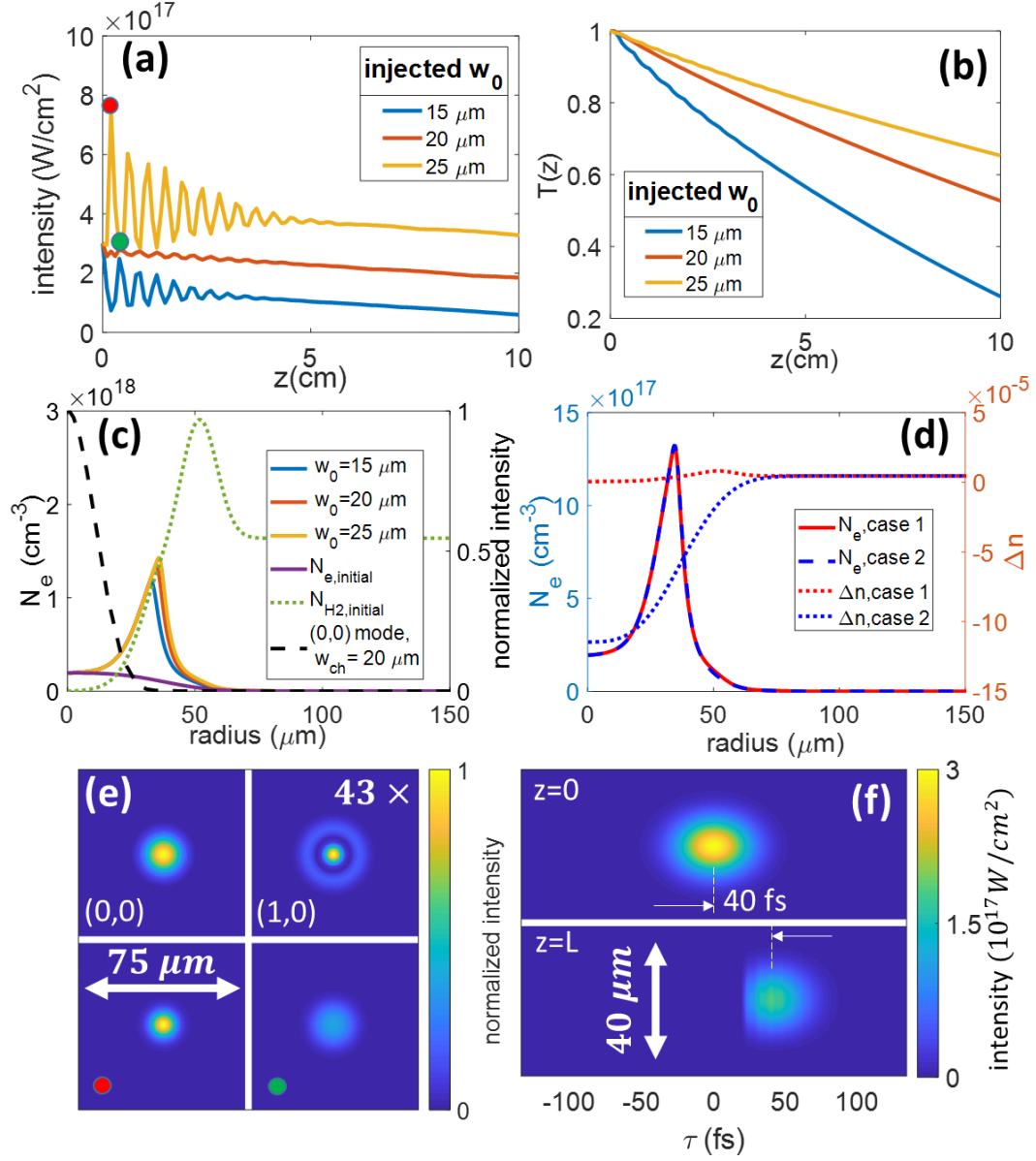


Figure 3.4. Peak input intensity at $z = 0$ is $I_0 = 3 \times 10^{17} W/cm^2$. (a) Peak laser intensity $I(r = 0, z, \tau_{peak})$ vs. z for various injection pulse waists w_0 . (b) z -dependent transmission efficiency, $T(z) = \epsilon_i^{-1} \int d\tau d^2r_{\perp} I(r_{\perp}, z, \tau)$, for the 3 input modes of (a), with variables defined in the text. (c) Similar $N_e(r)$ profiles generated by varying w_0 mode injection into the initial index profile shown (characterized by $N_{e,initial}$ and $N_{H2,initial}$). The dashed black curve is the computed ([83]), lowest order quasi-bound mode of the generated plasma waveguide. (d) Demonstration that a range of possible initial refractive index structures (plotted as refractive index difference $\Delta n = n - 1$) can support self-waveguiding, with almost identical plasma waveguides generated. The main feature necessary for self-waveguiding is an on-axis neutral gas density depression. Δn is the refractive index (e) Top 2 panels: (0,0) and (1,0) quasi-bound intensity modes of plasma waveguide in (b). The (1,0) mode

is shown scaled in intensity by $43 \times$. Bottom 2 panels: propagating mode at z locations of intensity maximum and minimum (green and red dots in (a)). (f) Pulse intensity envelopes in (r, τ) at the entrance ($z = 0$, top panel) and exit ($z = L = 10$ cm, bottom panel) of the index structure. The plots show pulse head erosion and group velocity delay in the self-generated waveguide. The pulse propagates right to left.

Figure 3.4(a) shows the peak pulse intensity $I(r = 0, z, \tau_{peak})$ as a function of distance along a $L = 10$ cm long structure (taken to be the measured 3.5 ns delay index profile from Fig. 3.2(a)) for varying w_0 for peak intensity $I_0 = 3 \times 10^{17}$ W/cm². Here $\tau = t - z/v_g$ is a time coordinate local to the pulse, v_g is the pulse group velocity, and τ_{peak} is the location of the pulse peak. Figure 3.4(b) plots the energy transmission efficiency $T(z) = \varepsilon_i^{-1} \int d\tau d^2\mathbf{r}_\perp I(\mathbf{r}_\perp, z, \tau)$ vs. z , where $\varepsilon_i = 0.53 I_0 \pi w_0^2 \tau$ is the pulse energy at the entrance to the index structure ($z = 0$) and the integral is over the pulse envelope and the index structure cross section. We will later discuss the physical origin of the oscillations and transmission decay in Figs. 3.4(a) and (b).

The structure cross section at $z = 2$ cm ($8z_0$) is shown in Fig 3.4(c) before and after injection of pulses with the varying waists from Fig. 3.4(a) and (b); the plasma waveguide profile ultimately generated is largely dictated by the initial neutral hydrogen profile (light green points) and *not* the injected beam waist. This is consistent with the measurements of Fig. 3.3(a). The computed lowest order quasi-bound mode for Fig. 3.4(c)'s similar plasma waveguide profiles is shown overlaid on the plot; it is near-Gaussian with a $1/e^2$ intensity waist $w_{ch} \sim 20 \mu\text{m}$, matching the $w_0 \sim 20 \mu\text{m}$ plot of Fig. 3.4(a), which has the least intensity oscillation. This suggests that injected pulses that are not matched to the ultimate waveguide profile could couple into the profile's higher order modes. Indeed, the source of the strong oscillation for the $w_0 \sim 15 \mu\text{m}$ and $w_0 \sim 25 \mu\text{m}$ cases in Fig. 3.4(a) is beating between the first two quasi-

bound modes excited when the injected pulse w_0 is not matched to the w_{ch} of the ultimate self-generated guide profile (Fig. 3.4(c)). These modes correspond to $(p, m) = (0, 0)$ and $(1, 0)$, where p and m are radial and azimuthal mode indices, and where $m = 0$ owing to negligible azimuthal asymmetry in the J_0 pulse-excited index structure. The intensity profiles of the $(0, 0)$ and $(1, 0)$ modes are shown in Fig. 3.4(e) (top two sub-panels). The mode intensity ratio is $\mathcal{R} = |a_{0,0}/a_{1,0}|^2 \sim 43$, where $a_{0,0}$ and $a_{1,0}$ are the modal expansion coefficients for the waveguide profiles of Fig. 3.4(c), and their computed wavenumber difference is $\Delta\beta = \beta_{0,0} - \beta_{1,0} = 2.3 \text{ cm}^{-1}$ [83]. Beating between these modes would have a period $\Lambda = 1/\Delta\beta = 4.3 \text{ mm}$ and a linear modulation visibility $= (I_{max} - I_{min})/(I_{max} + I_{min}) = 2\sqrt{\mathcal{R}}/(\mathcal{R} + 1) = 0.30$. The modulation period corresponds very closely to the oscillation period of the ‘non-matched’ curves of Fig. 3.4(a), while the modulation visibility is qualitatively consistent with what is really a transient nonlinear process involving ionization and plasma lensing, along with mode beating. The bottom two sub-panels of Fig. 3.4(e) show the intensity profiles at the maximum and minimum of the oscillations in Fig. 3.4(a) (red and green dots).

It is interesting to note in Fig. 3.4(a) that for the unmatched injection pulse with $w_0 = 25 \text{ }\mu\text{m}$, the peak intensity oscillations grow to as much as $\sim 2.7 I_0$ before they settle down to a steady level of $\sim 1.5 I_0$, followed by a smooth erosion decay to the channel exit. By contrast, the $w_0 = 15 \text{ }\mu\text{m}$ pulse oscillations never exceed I_0 and settle down to a steady level of $\sim 0.5 I_0$ before decaying smoothly to the exit. This is from injected energy being channeled into a spot w_{ch} , transversely compressing or expanding the pulse. This predicts an intensity enhancement

$(w_0/w_{ch})^2 \sim (25/20)^2 \sim 1.6 \times$ for the $w_0 = 25 \mu\text{m}$ pulse injection, and an intensity reduction $(w_0/w_{ch})^2 \sim (15/20)^2 \sim 0.6 \times$ for the $w_0 = 15 \mu\text{m}$ pulse injection. These estimates agree well with the simulations.

As discussed earlier, the on-the-fly generation of the plasma waveguide cladding is driven by the portion of the pulse's leading edge in excess of the hydrogen ionization threshold. Thus, earlier portions of the pulse below the ionization threshold will diffract away, and later portions, short of achieving full cladding confinement, will leak away until the cladding height and width is sufficient to stem further losses. Pulse head erosion is seen in Fig. 3.4(f), which compares (for the matched input $w_0 = 20 \mu\text{m}$) the input pulse at $z = 0$ and the pulse exiting the structure at $z = 10 \text{ cm}$. In addition, the exit pulse centroid delay ($\sim 40 \text{ fs}$) is consistent with the group velocity delay $\Delta\tau = L(v_g^{-1} - c^{-1}) \sim 45 \text{ fs}$, where the group velocity of the pulse in the plasma waveguide [99] is $v_g = (\partial\beta_{0,0}/\partial\omega)^{-1} = c[1 - N_{e0}/2N_{cr} - 2/(kw_{ch})^2]$. Here, $N_{e0} = 2 \times 10^{17} \text{ cm}^{-3}$ and $w_{ch} = 20 \mu\text{m}$ (from Fig. 3.4(b)), k is the laser vacuum wavenumber, and $N_{cr} = 1.7 \times 10^{21} \text{ cm}^{-3}$ is the critical density for $\lambda = 800 \text{ nm}$. We note that the self-generated waveguides shown thus far are tightly confining, with $L_{1/e} > 30 \text{ cm}$ for those simulated Fig. 3.4(c) and $L_{1/e} > 1.4 \text{ m}$ for the 88 mJ-generated guide of Fig. 3.3(a).

An important question arises: is there a class of initially prepared refractive index profiles that is best-suited to support self-waveguiding? Thus far in Fig. 3.4, we presented simulations corresponding to our experimental index profiles: a low density plasma core surrounded by an enhanced density neutral annulus. In Fig. 3.4(d) we simulate self-waveguiding in two other types of index profiles. In case 1, the initial

refractive index profile $\Delta n_{\text{case1}}(r)$ corresponds to a depressed on-axis neutral density surrounded by an annulus of enhanced density corresponding to the profile shown in Fig. 3.4(c). There is no plasma contribution. The other case, $\Delta n_{\text{case2}}(r)$, removes the enhanced neutral density annulus and keeps Fig. 3.4(c)'s central plasma core and neutral density depression. We note that the central concavity of both Δn_{case1} and Δn_{case2} profiles would ensure defocusing of a low intensity injected pulse, even more strongly for case 2. However, end-injection of the $I_0 = 3 \times 10^{17} \text{ W/cm}^2$, $w_0 = 20 \text{ }\mu\text{m}$ pulse into either the $\Delta n_{\text{case1}}(r)$ or $\Delta n_{\text{case2}}(r)$ structure yields the electron density profiles at $z = 2 \text{ cm}$ overlaid on the plot; they are nearly identical. Full propagation runs (not plotted here) show indistinguishable self-waveguiding; ionization occurs early enough that the remainder of each pulse sees the same plasma waveguide. We conclude that the most important feature of a prepared index profile is an adequately deep neutral density gas depression, where for hydrogen gas, the depth must exceed $\Delta N_{H2} > (2\pi r_e w_{ch}^2)^{-1}$ to confine a mode w_{ch} [83], where r_e is the classical electron radius. The presence of a neutral shock wall (as included in case 1) is unnecessary in general; however, under some circumstances it will reduce propagation leakage by providing higher and wider cladding. Plasma is required only to the extent that it contains the laser-provided thermal energy that drives the initial cylindrical blast expansion.

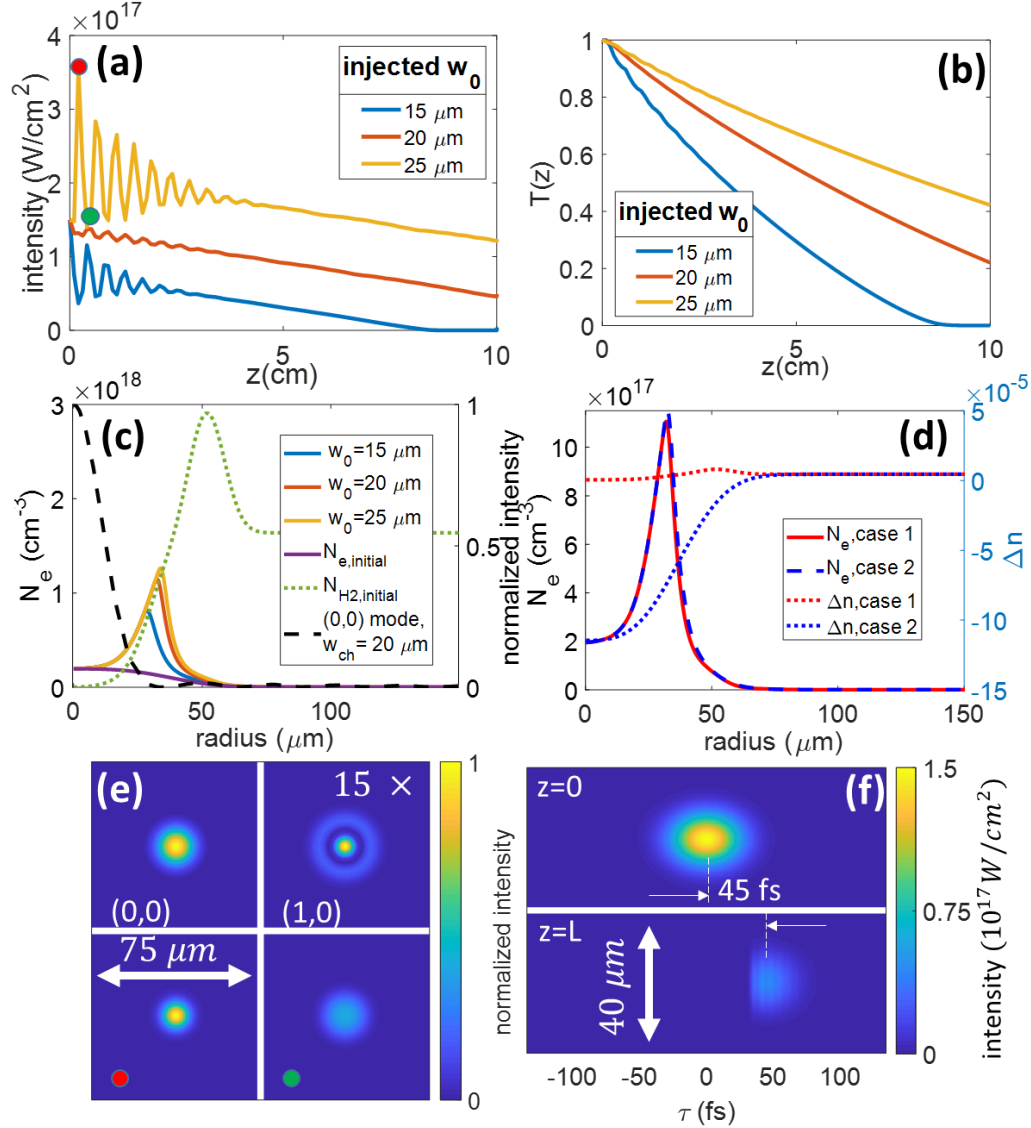


Figure 3.5. Peak input intensity at $z = 0$ is $I_0 = 1.5 \times 10^{17} \text{ W/cm}^2$. The panel descriptions are the same as in Fig. 4, except that in (e), the (1,0) mode is shown scaled in intensity by $15 \times$.

Figure 3.5 repeats the simulations of Fig. 3.4 except at a lower peak injected intensity, $I_0 = 1.5 \times 10^{17} \text{ W/cm}^2$. Here, the deleterious effects of a $2 \times$ intensity reduction are evident. The basic propagation features remain, such as the mode beating for unmatched injected beams, but we also observe the following: (i) The peak intensity $I(r = 0, z, \tau_{peak})$ and transmission $T(z)$ decline much faster with z (Fig. 3-

3.5(a) and 3.5(b)). (ii) While the electron profiles in 3.5(c) still follow the neutral density profile (light green points) the peak and width of the self-generated electron density cladding are smaller, especially for the $w_0 = 15 \mu m$ injected pulse. This accounts for the much faster leakage rate seen in Fig. 3.5(b)—in fact, the $w_0 = 15 \mu m$ pulse does not even survive to $z = 10$ cm. (iii) The intensity ratio of the (0,0) mode to the (1,0) mode is much smaller, $\mathcal{R} = |a_{0,0}/a_{1,0}|^2 \sim 15$ (for the $w_0 = 20 \mu m$ – induced electron density profile in Fig. 3.5(c)), giving larger linear modulation visibility $\nu = 0.48$. For all 3 profiles $\Delta\beta = \beta_{0,0} - \beta_{1,0} = 2.3 \text{ cm}^{-1}$ as before, giving rise to the same modulation periods seen in Figs. 3.4(a) and 3.5(a). The pulse energy redistribution effect is likewise seen in Fig. 3.5(a), with a peak intensity enhancement (after oscillations settle down) of $\sim 1.3 \times$ (for $w_0 = 25 \mu m$) and a $\sim 0.5 \times$ reduction (for $w_0 = 15 \mu m$), reduced effects owing to the weaker confinement of the self-generated guides. (iv) Pulse head erosion, as seen in Fig. 3.5(f), is more significant, as a larger portion of the leading edge is below the self-trapping threshold. This has the effect of moving the centroid of the exit pulse farther back.

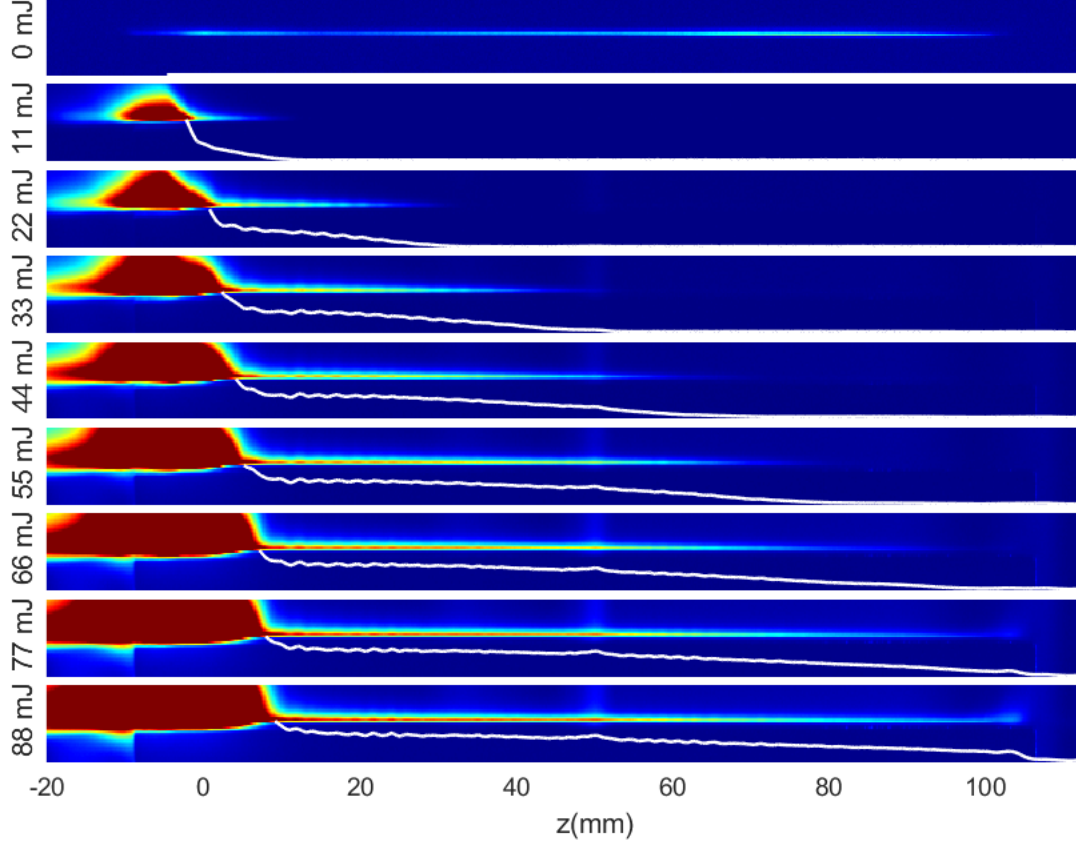


Figure 3.6. Full length hydrogen plasma fluorescence (recombination) images of self-waveguiding of $\lambda = 800$ nm, 40 fs pulses of increasing energy in a 10 cm long index structure generated in an elongated supersonic hydrogen jet by a 100 mJ J_0 pulse. $\Delta t_{injection} = 2.5$ ns into the profiles of Fig. 3.3(a). The pulse propagates left to right. Below each image is a central lineout (white curves). The decreasing fluorescence with z indicates pulse head erosion of the self-waveguiding pulse. The bump at $z \sim 50$ mm is from light scattering from an optical support post in the experimental chamber. The saturated (brown) image at the jet entrance is from the injected laser pulse overfilling the initial index structure. The slight bump at the channel exit is from ionization of neutrals by the exiting pulse. The top panel shows fluorescence induced by the J_0 pulse alone.

We now show results from our guiding experiments in the 10 cm long hydrogen jet. Figure 3.6 shows a sequence of hydrogen plasma recombination (fluorescence) images of the full jet for increasing injected pulse energy, for $\Delta t_{inj} = 2.5$ ns (into the 2.5 ns delay index profile of Fig. 3.2(a), resulting in the plasma profiles of Fig. 3.3(a)). The J_0 line focus is 0.8 mm above the nozzle's slot orifice, where the local atomic

hydrogen density is $N_0 = 3.2 \times 10^{18} \text{ cm}^{-3}$. The z -dependent fluorescence is proportional to the local radial integral of the electron density, $F(z) \propto \int_0^{r_{\max}} dr 2\pi r N_e(r, z)$, and tracks the z -dependent free electron charge generated by the guided pulse; it is a qualitative measure of the axial extent of guiding.

As the injected laser energy increases, the pulse propagates progressively longer distances because a smaller fraction of its leading edge is below the thresholds for hydrogen ionization and self-trapping. That fraction continually erodes from the front of the pulse, reducing its energy until a self-trapping structure can no longer be generated and the remaining pulse diffracts away. We see that for injected energy $> 75 \text{ mJ}$, the pulse makes it to the end of the 10 cm long structure and then expands, where it excites a small burst of fluorescence in residual neutral gas at the exit. The much larger fluorescence signal at the entrance (shown saturated) may be from ionization of the index structure by the wings of the injected beam mode. The oscillations seen in the simulations of Figs. 3.4(a) and 3.5(a) may contribute to these wider wings, as seen in Figs. 3.4(e) and 3.5(e). Below each fluorescence image is a central lineout (white curves), from which an erosion rate of $\sim 6 \text{ mJ/cm}$ is estimated from the downward slope. This is mainly a diffraction and leakage loss: absorption is negligible, as the laser energy expended on ionization is $< \sim 1 \text{ mJ}$ for all guides measured in our experiments. Measurable end mode transmission occurred only when the fluorescence showed ‘burn through’ to the end of the jet. The improved transmission with increasing pulse energy is also seen in the simulations by comparing $T(z)$ in Figs. 3.4(b) and 3.5(b). These curves show an energy erosion rate between 7.5 mJ/cm (Fig. 3.4, $I_0 =$

$3 \times 10^{17} \text{ W/cm}^2$, $w_0 = 25 \mu\text{m}$) and 4 mJ/cm (Fig. 5, $I_0 = 1.5 \times 10^{17} \text{ W/cm}^2$, $w_0 = 25 \mu\text{m}$), consistent with our measured erosion rate of $\sim 6 \text{ mJ/cm}$.

A feature common to the fluorescence images for injected pulse energy $> 25 \text{ mJ}$ is a prominent axial modulation. In Fig. 3.7, we plot central lineouts from fluorescence plots for injected 88 mJ pulses at a range of delays Δt_{inj} , with their z-averaged modulation periods Λ indicated on the plot. Each of these delays corresponds to an index profile plotted in Fig. 3.3(d). Based on the simulations and analysis of Figs. 3.4 and 3.5, we attribute the modulations to beating between the (0,0) and (1,0) modes. Using the profiles of Fig. 3.3(d), we calculate the quasi-bound mode wavenumber spectrum [99] and determine $\Lambda_{sim} = (\Delta\beta)^{-1}$, where $\Delta\beta = \beta_{0,0} - \beta_{1,0}$. These values are added to the panels in Fig. 3.7 and show mostly good agreement with the measured z-averaged periods Λ (except for the $\Delta t_{inj} = 5 \text{ ns}$ case, for reasons unclear).

For successful self-waveguiding of intense laser pulses, the self-generated guide must be able to support at least one low loss quasi-bound mode, as in Figs. 3.4(c) and 3.5(c). An approximate condition on the waveguide profile for trapping and guiding of a (p, m) mode is

$$\Delta N_e > (2p + m + 1)^2 / \pi r_e w_{ch}^2 \quad (3.1)$$

where $\Delta N_e \approx N_e(r = w_{ch}) - N_e(r = 0)$ and w_{ch} is the radius of the corresponding fundamental mode [99]. The requirement of minimum ΔN_e implies a threshold initial hydrogen gas density for guiding. This is seen in Fig. 3.8(a) (for injection pulse energy 88 mJ and $\Delta t_{inj} = 2.5 \text{ ns}$), which shows a sharp onset of guiding over the full 10 cm channel for $N_{H_2} > 10^{18} \text{ cm}^{-3}$. Shown are plots of pulse energy exiting the guide and transmission efficiency (pulse energy at guide exit / energy at entrance), along with exit

modes measured for low, intermediate, and higher hydrogen densities. At the onset of guiding, only the (0,0) mode exits the guide. As H_2 density increases, the contribution of the (1,0) mode becomes more significant, consistent with Eq. (3.1), which implies that density should increase by $\sim 4 \times$ for $p = 1$ mode confinement. The presence of the (1,0) mode is consistent with the mode beating seen in Fig. 3.6. Injected pulse transmission vs. energy is shown in Fig. 3.8(b) (for $\Delta t_{inj} = 2.5$ ns and $N_{H_2} = 3 \times 10^{18} \text{ cm}^{-3}$), where the low end of the energy axis is the minimum needed for burn-through of the 10 cm long structure. The increased transmission at high energy (approaching $\sim 10\%$) under these conditions reflects the reduced fractional erosion of the head of the pulse. The highest transmitted energy of 23 mJ corresponds to a peak self-waveguided exit intensity of $1.5 \times 10^{17} \text{ W/cm}^2$, in the weakly relativistic regime (normalized vector potential $a_0 \sim 0.3$). Upstream from the exit were even higher intensities, as is clear from Figs. 3.4(a,b) and 3.5(a,b).

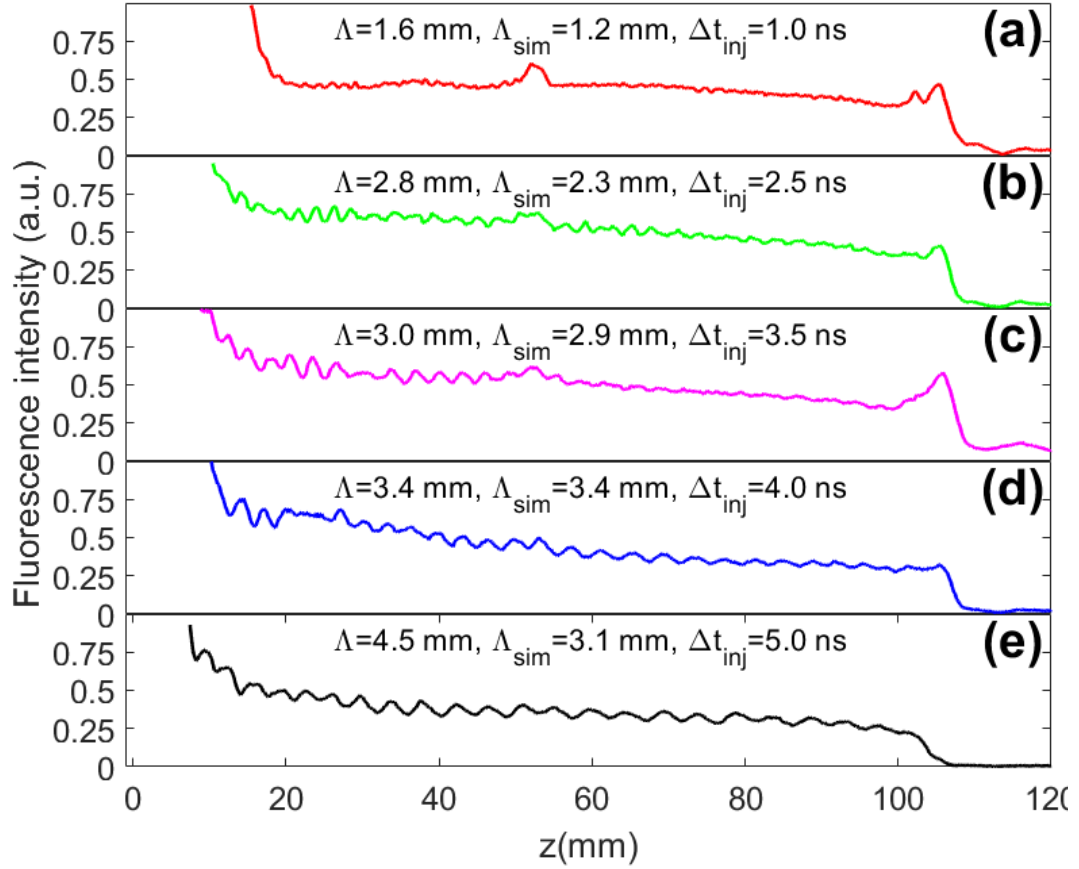


Figure 3.7. Full length hydrogen plasma fluorescence (recombination) image lineouts of self-waveguiding (of $\lambda = 800$ nm, $w_0 = 25$ μ m, 40 fs, 88 mJ pulses) in a 10 cm long index structure generated in an elongated supersonic hydrogen jet by a 100 mJ J_0 pulse. Here the injection delay Δt_{inj} is varied from 1 ns to 5 ns, slightly changing the self-generated plasma waveguide (as seen in Fig. 3.3(d)) and its quasi-bound modes and wavenumbers $\beta_{0,0}$ and $\beta_{1,0}$, leading to a variation in the mode beating period $\Lambda_{sim} = (\Delta\beta)^{-1}$, where $\Delta\beta = \beta_{0,0} - \beta_{1,0}$. The simulated period Λ_{sim} and measured profile-average period Λ are in reasonable agreement. The bump at $z \sim 50$ mm is from light scattering from an optical support post in the experimental chamber.

The simulations in Figs. 3.4 and 3.5 show a strong sensitivity of $T(z = L)$ to I_0 and w_0 , which is likely reflected in our transmission measurements. Given the uncorrected astigmatism from the 0.5 m spherical focusing mirror, we attribute the mismatch between the transmission measurements and the simulations of Figs. 3.4(b) and 3.5(b) to the uncorrected mode. Nevertheless, our results are reasonably consistent with the

simulations. Importantly, it is clear that higher intensity injected pulses will experience much higher transmission fractions.

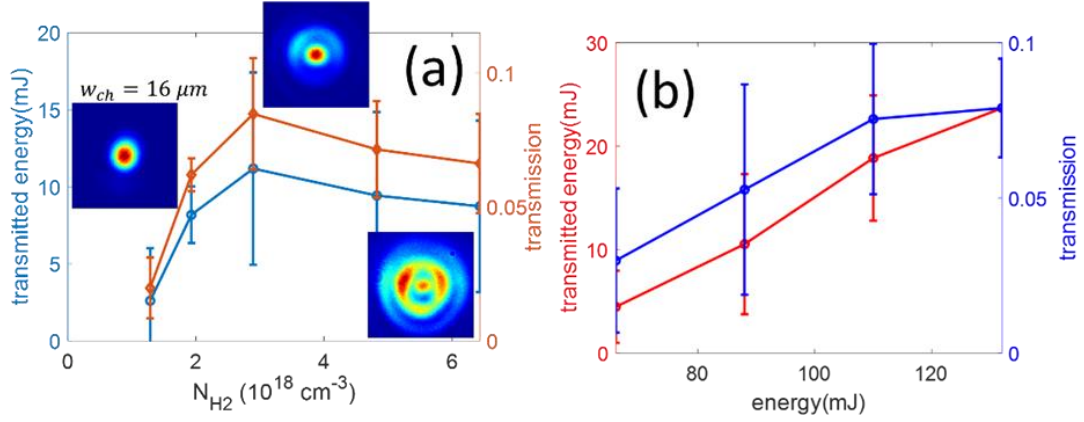


Figure 3.8. (a) Plot showing hydrogen density threshold for self-waveguiding (injection pulse energy 88 mJ, $\Delta t_{inj} = 2.5$ ns). As density increases, both (0,0) and (1,0) modes are seen at the channel exit, with the contribution of the (1,0) mode increasing. The slightly reduced transmission at higher density may be from leakage of the (1,0) mode. (b) Injected pulse transmission vs. energy (for $\Delta t_{inj} = 2.5$ ns and $N_{H_2} = 3 \times 10^{18} \text{ cm}^{-3}$), where the low end of the energy axis is the minimum needed for burn-through of the 10 cm long structure. Peak self-waveguided intensity exiting the guide is $1.5 \times 10^{17} \text{ W/cm}^2$.

The relatively low transmission efficiency measured here for pulses in the $\sim 10^{17} \text{ W/cm}^2$ range should not be considered as an argument against implementing the self-waveguiding method. Rather, it could be considered the cost of generating a plasma waveguide for a follow-on pulse of much higher intensity. Such a follow-on pulse can be used, for example, for LWFA. To illustrate this operating scenario, we performed a 3-pulse experiment in our 10 cm gas jet: a J_0 pulse ($\lambda = 800 \text{ nm}$) generated the initial refractive index structure, a pump pulse ($\lambda = 800 \text{ nm}$) injected into the structure at $\Delta t_{inj} = 2.5$ ns generated the plasma waveguide, and a probe pulse ($\lambda = 400 \text{ nm}$) followed the waveguide generation pulse at delay ~ 5 ps. The probe pulse mode was not optimized and overfilled the waveguide entrance. Figure 9 is a plot of probe

pulse transmission as a function of pump pulse energy. Above the curve are shown $\lambda=400$ nm probe output mode images and below the curve are the corresponding self-waveguiding exit mode images at several injected pump pulse energies. Most notable is that all modes, for both pump and probe at $\lambda = 800\text{nm}$ and $\lambda = 400\text{nm}$, are similar, with $w_{ch} = 16 - 18 \mu\text{m}$. This is strong evidence of two important features: First, it clearly reemphasizes that the prime determiner of the plasma waveguide structure left by a self-waveguiding pulse is the initial neutral density profile induced by the index-structuring J_0 pulse. And second, it verifies that the mode structure of a given plasma waveguide is wavelength-independent [99].

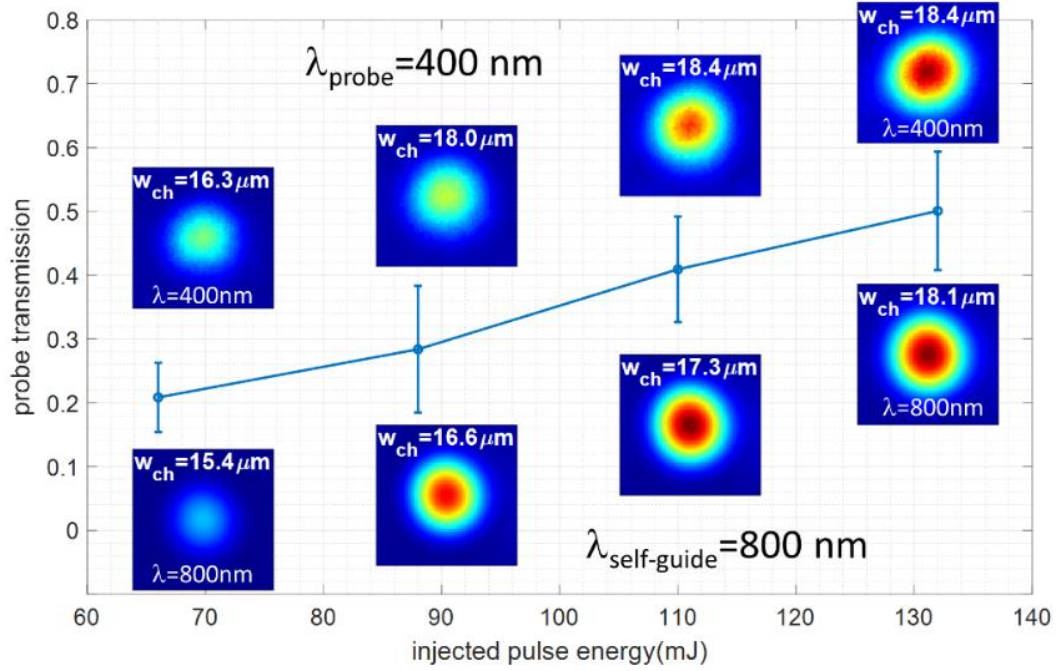


Figure 3.9. Pump-probe experiment, where the pump is the self-waveguiding pulse ($\lambda=800$ nm) injected at $\Delta t_{inj} = 2.5$ ns after the 100 mJ J_0 pulse, followed at 5 ps delay by a probe pulse ($\lambda=400$ nm). The corresponding initial index structure is plotted in Fig. 2(a), and its generated waveguide is shown in Fig. 3(a). The self-waveguided $\lambda=800$ nm exit modes are shown below the curve, with the corresponding $\lambda=400$ nm probe exit modes shown above the curve. The mode invariance ($w_{ch} = 16 - 18$ μm) illustrates that (i) the resulting plasma waveguide structure is nearly independent of the energy of the self-waveguiding pulse and (ii) the plasma waveguide mode structure is wavelength independent.

3.4 Comparison of the Two Methods

In chapters 2 and 3, we have presented two different methods of forming plasma waveguides with optical field ionization: self-waveguiding and the two Bessel method. In this section, we will compare the two methods on the basis of energy efficiency.

We begin with computing the cost per unit length for the two-Bessel pulse method [11] using a J_q Bessel beam with $\alpha = 3^\circ$, $\tau_{FWHM} = 50$ fs and a super Gaussian

profile. We chose a 10 mm FWHM beam diameter and the order $q = 18$ to match the length and width of the self-waveguiding-generated plasma. As is shown in chapter 2, the Bessel beam profile is unaffected by any preexisting plasma and the atomic density was assumed to be equal peak electron density in the self-waveguiding profile (blue curve in Fig. 3.10). The ionization rate was then calculated using a modified ADK model [100] and the z-averaged plasma profiles shown in Fig. 3.10 were found.

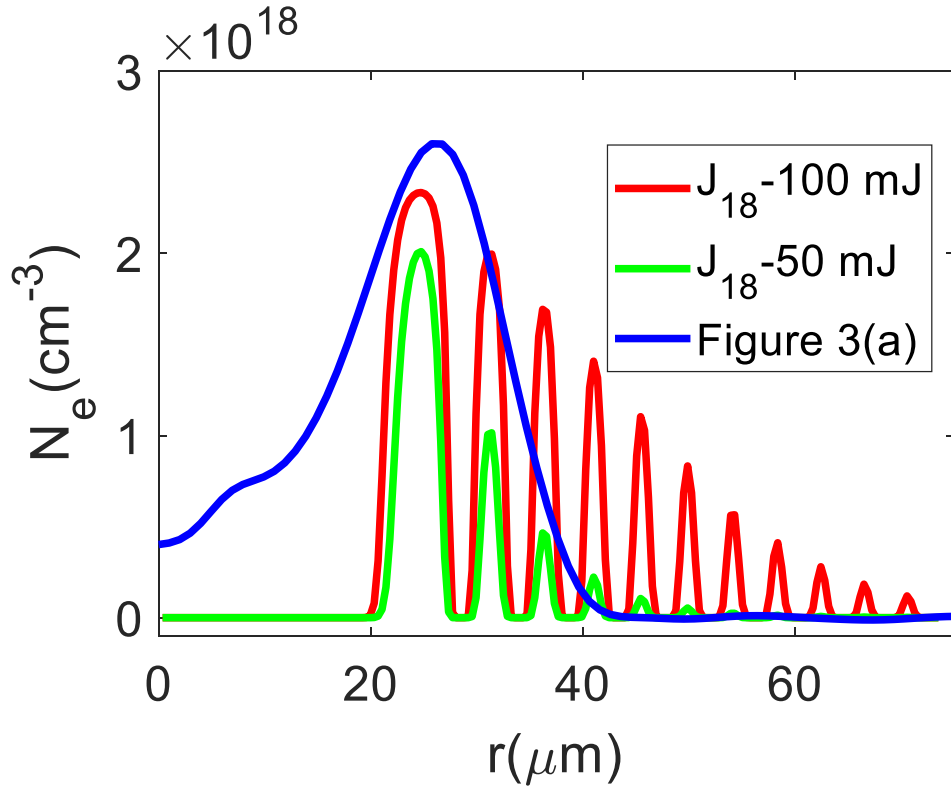


Figure 3.10. Energy requirement comparison. Blue curve: Plasma density profile shown in Fig. 3(a) of the main text. Red line: Plasma density profile generated by a J_{18} Bessel beam pulse with 100 mJ energy, assuming a background hydrogen atom density $N_H = 2.5 \times 10^{18} \text{ cm}^{-3}$. Green line: Plasma density profile generated by a J_{18} Bessel beam with 50 mJ energy with the same background N_H .

For a Bessel beam of order q , the intensity distribution can be written (for optical axis approach angle $\gamma = 2\alpha \ll 1$) as $I(r, z) = I(\rho_\gamma)(2\pi k \rho_\gamma \sin \gamma) J_q^2(kr \sin \gamma)$, where $\rho_\gamma = z \tan \gamma$ is radial position on the beam. If the plasma cladding is generated by the

first ring of J_q , then we require $kr_{max}\sin\gamma = \xi_{q1}$, where r_{max} is the peak location of the first ring, where $J_q'(\xi_{q1})=0$. The z-dependent peak intensity along the first ring is $I(r_{max}, z) = I(\rho_\gamma)(2\pi k\rho_\gamma \sin\gamma)J_q^2(\xi_{q1})$. Using the length of the focal line $L = R/\tan\gamma$ for a beam of radius R , the average ring intensity along the focal line is

$$\langle I(z) \rangle = \frac{1}{L} \int dz I(r_{max}, z) = \frac{k\sin\gamma}{R} J_q^2(\xi_{q1}) \int_0^R 2\pi\rho_\gamma I(\rho_\gamma) d\rho_\gamma \quad (3.2)$$

Using $\langle I(\rho) \rangle = (\pi R^2)^{-1} \int_0^R 2\pi\rho_\gamma I(\rho_\gamma) d\rho_\gamma$, we get $\langle I(z) \rangle = k\pi R\sin\gamma J_q^2(\xi_{q1}) \langle I(\rho) \rangle$.

Thus, to supply a focal line average threshold ring intensity $\langle I(z) \rangle_{th}$ to ionize hydrogen, the energy in the input beam (of pulse duration τ) must be

$$\varepsilon \approx \langle I(\rho) \rangle \pi R^2 \tau = \langle I(z) \rangle_{th} \frac{R\tau}{k\sin\gamma J_q^2(\xi_{q1})} \approx \langle I(z) \rangle_{th} \tau \frac{L}{kJ_q^2(\xi_{q1})} \quad (3.3)$$

The energy required per unit length is therefore

$$\frac{\varepsilon}{L} \approx \frac{\langle I(z) \rangle_{th} \tau}{kJ_q^2(\xi_{q1})} \quad (3.4)$$

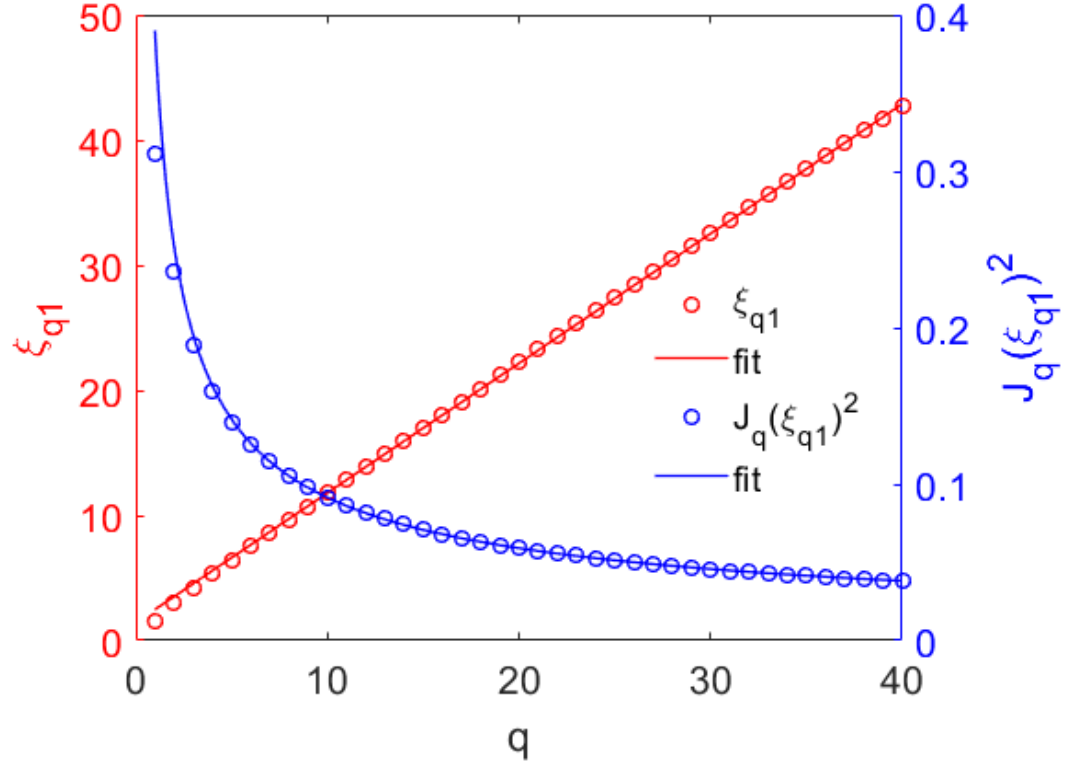


Figure 3.11. Computation of ξ_{q1} and $J_q^2(\xi_{q1})$ vs. q (red circles and blue circles). Fits are overlaid: $(\xi_{q1})_{fit} = 1.03q + 1.44$ and $(J_q^2(\xi_{q1}))_{fit} = 0.39q^{-0.63}$.

Figure 3.11 plots ξ_{q1} and $J_q^2(\xi_{q1})$ vs. q , with the fits $(\xi_{q1})_{fit} = 1.03q + 1.44$ and $(J_q^2(\xi_{q1}))_{fit} = 0.39q^{-0.63}$ overlaid. Taking $\xi_{q1} \approx q$ and using $w_{ch} = a\xi_{q1}$ where $a \lesssim 1$ gives $w_{ch} = qa/k\sin\gamma$ or $q = kw_{ch}\sin\gamma/a$. Therefore, we can write $(J_q^2(\xi_{q1}))_{fit} = 0.39q^{-0.63} = 0.39(kw_{ch}\sin\gamma/a)^{-0.63}$, giving $J_q^{-2}(\xi_{q1}) \propto w_{ch}^{0.63}$. Therefore, the energy cost per unit length is

$$\frac{\varepsilon}{L} \propto w_{ch}^{0.63} \quad (3.5)$$

increasing sub-linearly with w_{ch} .

To generate the plasma waveguide profile of Fig. 3.3(a) (for which $w_{ch} = 20 \mu\text{m}$) over 10 cm using the two-Bessel method, this method yields an energy requirement of $\sim 5 - 10 \text{ mJ/cm}$, in the same range as the $\sim 6 \text{ mJ/cm}$ estimated for the

self-waveguiding pulses of the current paper. However, as the required w_{ch} increases further, the energy required in the self-waveguiding pulse must scale as $\sim w_{ch}^2$ to maintain sufficient intensity in the wings to ionize the neutral gas, which suggests less efficiency at larger channel sizes. Larger channel sizes will be required for optimal LWFA, since geometric effects reduce the group velocity of the laser for small channel sizes, thus limiting the dephasing length and energy gain of the accelerator [14].

This comparison is also only valid for the intensity range discussed in this chapter, with $a_0 < 1$. At the intensities required for electron acceleration, $a_0 > 1$, the losses in the self-waveguiding process will be different. Self-waveguiding at those intensities and acceleration will be discussed in the following chapter.

It is also useful to consider the tunability of the two methods. In the self-waveguiding the method, the mode structure is determined by the combined neutral gas and plasma profile at the time delay of the self-waveguided pulse. This profile is in turn determined by the size and temperature of the initial J_0 -generated plasma column. In the two Bessel method, the mode size can be chosen more independently by the choice of higher order Bessel beam. However, it is possible that at high guided beam intensities, further ionization by the wings of the guided beam in the two-Bessel case will end up producing the same waveguide as self-waveguiding. Further experiments will be needed to address this question.

Chapter 4: GeV Electron Beams from a Self-Waveguided Laser Wakefield Accelerator

4.1 Introduction

The previous two chapters introduced new techniques for creating plasma waveguides and demonstrated these techniques at guided beams of $a_0 \sim 0.1$, an order of magnitude too low for LWFA. In this chapter, we apply the self-waveguiding technique (Chapter 3) in the first demonstration of GeV electron acceleration in an optically generated plasma waveguide and use PIC simulations to analyze the results. A laser pulse with $a_0 = 1.5$ was self-waveguided in low density hydrogen plasma over 20 cm, driving a wake and accelerating electrons up beyond a GeV. In the first stage of the self-waveguiding process, an index-structuring Bessel beam generates a neutral density depression in a long hydrogen gas jet. The early leading edge of the high intensity pulse injected into this neutral depression generates a waveguide via optical field ionization. The waveguide then confines the high intensity portion of the pulse envelope, which drives a plasma wave that can accelerate electrons. An Argon-doped 5 mm section of the gas jet provides a source for localized ionization injection of electrons.

Laser wakefield acceleration (LWFA) has been the subject of continuous experimentation and improvement since it was first proposed in 1979 [13]. In recent years, LWFA experiments using petawatt-scale laser systems have achieved electron energies up to 8 GeV [65], with future designs aiming at 1 TeV using 100 stages, each of 10 GeV [17,52]

The maximum accelerating field of the wake E_z and the energy gain, $\Delta\mathcal{E} = qE_z L_{acc}$, are limited by laser power and electron density. Here q is the electron charge and L_{acc} is the acceleration length. To maximize L_{acc} there are two requirements. First, the electron density must be sufficiently low that $L_{acc} > L_{depl}$, L_{deph} , the density-dependent depletion and dephasing lengths described in section 1.5. Second, some sort of optical guiding is needed so that E_z is maintained over distances much greater than the laser's diffraction length. One approach is to take advantage of relativistic self-guiding in the high a_0 nonlinear bubble regime, but this regime limits the flexibility of the accelerator and makes it difficult to optimize the emittance of the electron beam [52].

Another approach is to operate the accelerator in the quasi-linear regime, with $a_0 \sim 1$, and use a plasma waveguide, which were introduced in Chapter 1.6, to guide the drive laser and increase the acceleration length [47,61,101,102]. Analogously to optical fibers, plasma waveguides consist of a lower plasma density, higher index of refraction “core” and a higher plasma density, lower index “cladding”. The earliest plasma waveguides were created by the hydrodynamic cylindrical shock expansion of a plasma. In this scheme, a ~ 100 ps axicon generated Bessel beam ionizes a column of plasma through collisional and avalanche ionization, and then the column is further heated to ~ 100 eV by inverse bremsstrahlung (IB) heating [58]. The hot plasma then expands at the ion sound speed, forming shock walls that serve as the cladding of the waveguide. However, efficient IB heating requires electron densities $> 5 \times 10^{18} \text{ cm}^{-3}$ or the use of a highly absorptive media such as clustered gas [60]. The higher plasma

density in these cases shorten the dephasing and pump depletion lengths, limiting possible energy gains to a few hundred MeV [61].

The guiding approach used in previous experiments to achieve GeV level acceleration is the capillary discharge waveguide [65,102,103]. In this scheme, an electric discharge ionizes and heats the gas inside a material capillary (usually ceramic or sapphire). The steady state plasma pressure profile—uniform across the capillary cross section—ensures that the plasma density is lowest at the center (where the plasma temperature is highest), and highest near the wall (where the temperature is lowest) [62,63]. Because of finite capillary diameter and the need for a stable and uniform discharge, the resulting plasma waveguide tends to be wide and of higher density than is optimal for electron acceleration. This limitation has been mitigated somewhat in recent LWFA experiments using an auxiliary 8-ns, 300 mJ heater beam to further deepen the channel [64,65]. Discharge capillaries are also susceptible to a high heat load and material erosion, especially at high repetition rates; high power laser damage to the solid capillary is also an issue [67].

Bypassing the limitations of IB-heated plasma waveguides and capillary waveguides, recent experiments have demonstrated the formation of plasma waveguides using ultrashort Bessel beam pulses for initial generation of the plasma column. Here, an ultrashort ‘index structuring’ J_0 Bessel beam pulse generates a plasma column in a neutral gas through optical field ionization (OFI) which then radially expands and drives a shock wave on the periphery of the plasma. The result, when ionized, is the desired plasma density profile for a waveguide: a reduced central plasma density, increasing radially to the shock. Although ionization and heating per

particle of OFI-generated plasma is independent of density, it is limited to the ~ 10 eV [104], the ponderomotive potential at the $\sim 10^{14}$ W/cm² ionization threshold of hydrogen. This limits the height of the shock, limiting the waveguide cladding to prevent guided mode leakage—as a result, guiding over more than a few mm is difficult [94,105].

Importantly, the expanding shock at the plasma boundary encompasses all ionization stages including the neutrals, which dominate on the leading edge of the shock. This ‘bump’ in neutral density can be ionized by a second, higher order ultrashort Bessel beam (Chapter 2, [106]), or by the leading edge radial wings of an injected intense pulse (Chapter 3, [107–109]). In this Chapter, the self-waveguiding approach of Chapter 3 is applied with a much higher energy guided pulse, from Colorado State University’s ALEPH laser [110]. The higher energy pulse enables the investigation of self-waveguiding at sufficiently high intensities to accelerate electrons and waveguide is extended to 20 cm. Using a sectioned gas jet, electrons are injected via ionization injected electrons accelerated beyond 1.5 GeV.

4.2 Experimental Setup

The experimental setup is detailed in Figure 4.1. The gas jet is segmented, with a front 0.5 cm section for electron injection followed by a 19.5 cm acceleration section. The injection section gas is a 95% hydrogen, 5% argon mix for ionization injection [50], while the acceleration section gas is hydrogen. A 8.5 J, 45 fs, $\lambda = 800$ nm waveguiding pulse was focused by a f/25 off-axis paraboloid (OAP) to a $w_0 = 35$ μ m focal spot, yielding a peak normalized vector potential $a_0 = 1.5$. This corresponds to a peak vacuum intensity of 4.8×10^{18} W/cm². The pulse is of

sufficiently high intensity that it generates its self-confining waveguide cladding, as described above, ionizes argon in the ionization injection section, and drives a plasma wave to accelerate the injected electrons along the 19.5 cm acceleration section. The initial plasma column was formed by a 1 J, $\lambda = 800$ nm, linearly polarized index structuring pulse, which was focused by a transmissive ring grating to form a J_0 Bessel beam with approach angle $\alpha = 2.3$. The Bessel beam was sufficiently intense to fully ionize the peak of a plasma column with FWHM=18 μm along the entire length of the gas jet.

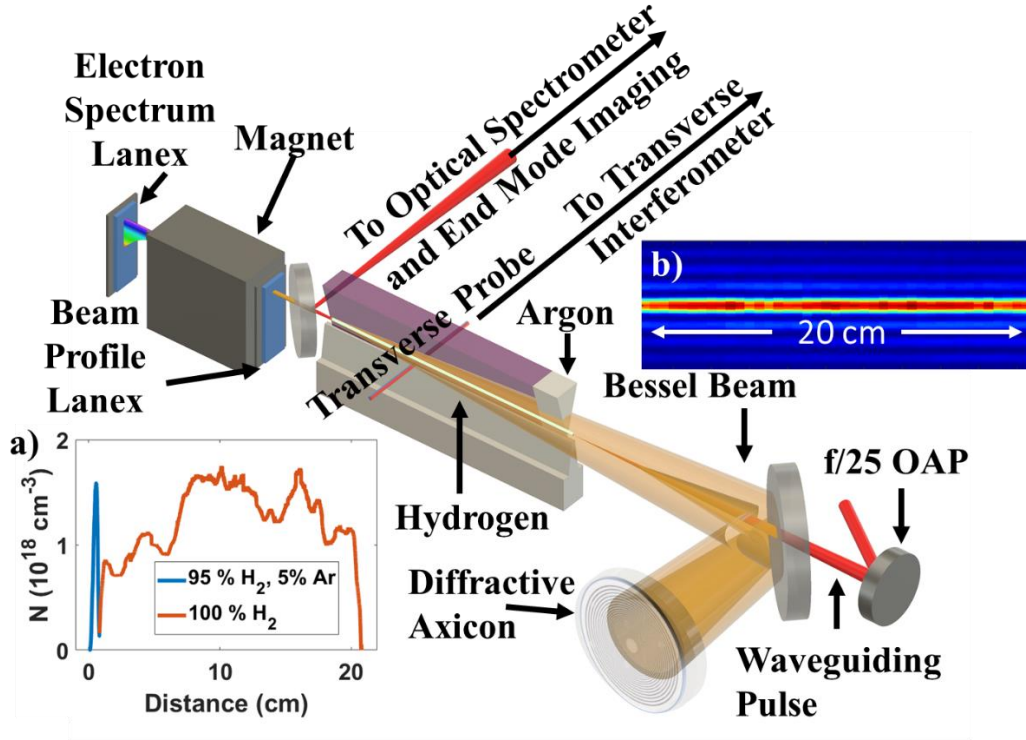


Figure 4.1. Experimental Setup. An index-structuring pulse (1 J , $\lambda_0 = 800 \text{ nm}$, linearly polarized) is focused by a transmissive ring grating (groove spacing $d = 20 \mu\text{m}$) forming a J_0 Bessel beam with approach angle $\alpha = 2.3^\circ$. The Bessel beam axial profile is shown in (a). The input beam is reflected off a mirror with a hole in it and forms a Bessel beam over a two section, 20 cm, supersonic gas nozzle, with first 0.5 cm nozzle section puffing a 5% Ar/95% H_2 mix gas jet and the following 19.5 cm section puffing a pure H_2 jet. A 8.5 J, 45 fs, $\lambda = 800 \text{ nm}$, linearly polarized pulse is focused by a f/25 off-axis paraboloid through the hole in the mirror to a $w_0 = 35 \mu\text{m}$ focal spot at the beginning of the jet, and is then self-waveguided in the J_0 beam-generated gas depression. The exit end of the guide is imaged by an exit mode imager and the transmitted optical spectrum is measured by an Ocean Optics spectrometer HR 2000+. Electron beam profiles and energy are measured by LANEX fluorescent screens and a permanent magnet spectrometer. A two-color ($\lambda = 400$ and 800 nm) interferometric probe is directed transversely through the index structure and into a folded wave-front interferometer. (b) shows a lineout of the neutral gas density profile of the gas jet, measured with plasma recombination fluorescence for the second section, after optical field ionization from the J_0 pulse.

4.3 Guiding Results

Figures 4.2-4.4 show the results of initial guiding experiments taken with initial atomic hydrogen density in the range $N_0 = 6 - 16 \times 10^{17} \text{ cm}^{-3}$ and peak vector potential in the range $a_0 = 0.6 - 1.5$. In these experiments, only the second 19.5 cm long section of the gas jet, with pure hydrogen, was used. The waveguide was formed by the self-waveguiding process described in detail in Ch. 3 and ref. [108], with the main difference being the intensity of the guided pulse. Since the energy of electrons liberated by OFI—approximately equal to the ponderomotive potential at the instant of ionization (here ~ 10 eV)-- is purely dependent on the atomic species, the plasma column generated by the J_0 beam expands at a rate independent of initial gas density N_0 . At the same time, 2-color interferometry data (Fig 3.2) shows that the on-axis density electron density drops to $\sim 0.16 N_0$.

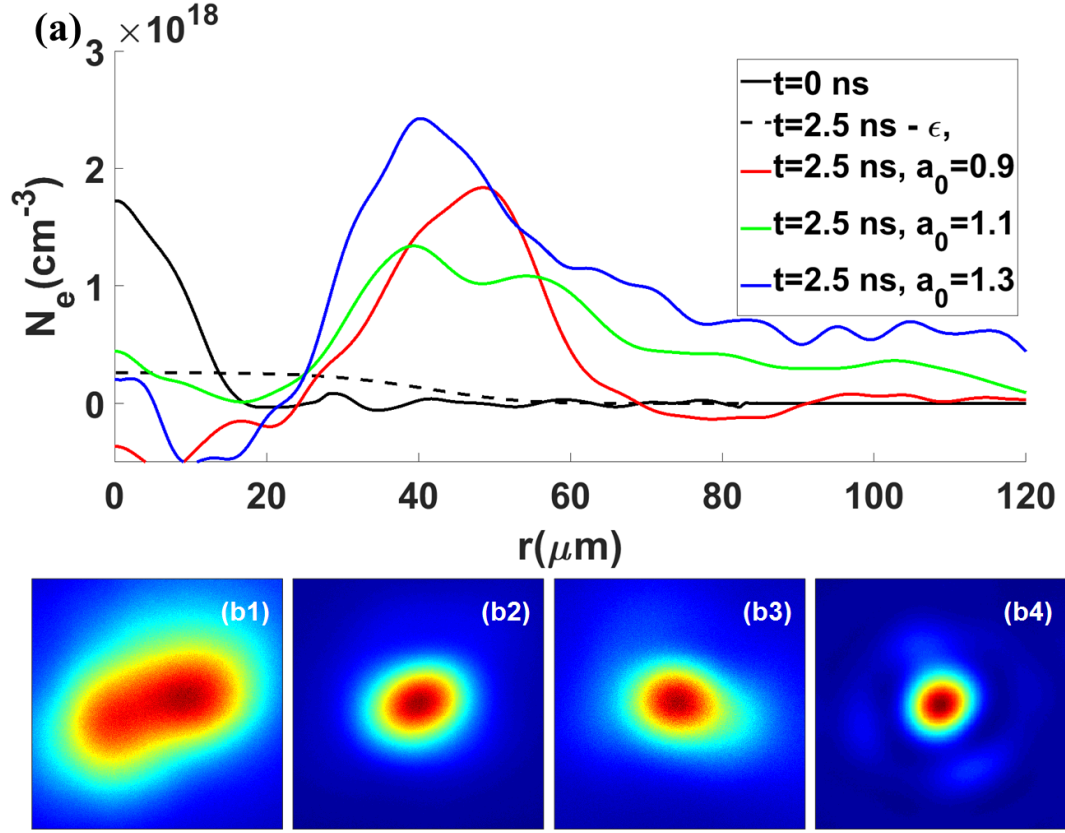


Figure 4.2 (a) Plasma density as indicated by the legends. The black dotted line is the plasma density immediately prior to the arrival of the self-waveguided pulse, measured by two-color interferometry (Fig 3.2) and scaled with respect to the initial plasma density. (b1)-(b4) Transmitted beam profile with increasing laser energy, corresponding to the three colored lines in (a). (b4) is the $f/25$ focal spot. All the four panels are normalized to its maximum respectively. The window size is $200 \mu\text{m} \times 200 \mu\text{m}$.

The leading edge of the self-waveguiding pulse then ionizes the neutral density profile, forming a plasma waveguide, the profile of which is shown in Fig. 4.2(a) for a range of a_0 at the vacuum focus of the self-waveguiding pulse. The peak intensity of the self-waveguided pulse is high enough that ionization occurs over radial distances more than $75 \mu\text{m}$ from the axis, forming the large plasma column shown in the Figure. The unphysical negative density values near the axis come from asymmetries and noise in the extracted phase profile, exacerbated by the Abel

Inversion. Figure 3.2 shows that at least 70% of the initial charge of the plasma is conserved for over 4 ns, so the on-axis density can be assumed to less than 130% of the density immediately before the waveguiding pulse, which is plotted at the black dotted line in Fig. 4.2 (a). The actual energy required to ionize the waveguide is <10 mJ for the entire 20 cm. The total waveguiding cost, including propagation losses, was estimated to be ~ 7 mJ/cm based on the measurements in Chapter 3 with a smaller mode size. This would correspond to 34 mJ/cm or 669 mJ for the whole 20 cm guide, assuming a scaling of w_{ch}^2 .

Figure 4.3 shows the transmission efficiency for various plasma densities and vacuum a_0 of the waveguiding beam. The transmission efficiency was measured by comparing the number of counts on the end mode imaging camera with and without a waveguide. These measurements should be taken to represent a lower bound on the actual transmission, as the measurements did not account for redshifting of the pulse due to the excitation of plasma waves. A sample optical spectrum showing the redshift centered around $1.1 \mu m$ is shown in figure 4.3 (b), but the end mode camera was not sensitive beyond $1 \mu m$ and the uncertainty in the spectral amplitude is large at this wavelength, so a corrected efficiency cannot reliably be estimated. The redshift could explain the apparent maximum of transmission at $a_0 = 0.8$ shown in figure 4.3 (a), as the redshift is expected to increase with laser energy [53]. Figure 4.3 (b) also shows the transmitted optical spectra found in a PIC simulation using the code FBPIC [33]. The simulations show that for the highest energies guided, over 70% of the guided energy is out of the detectable spectral range of the camera.

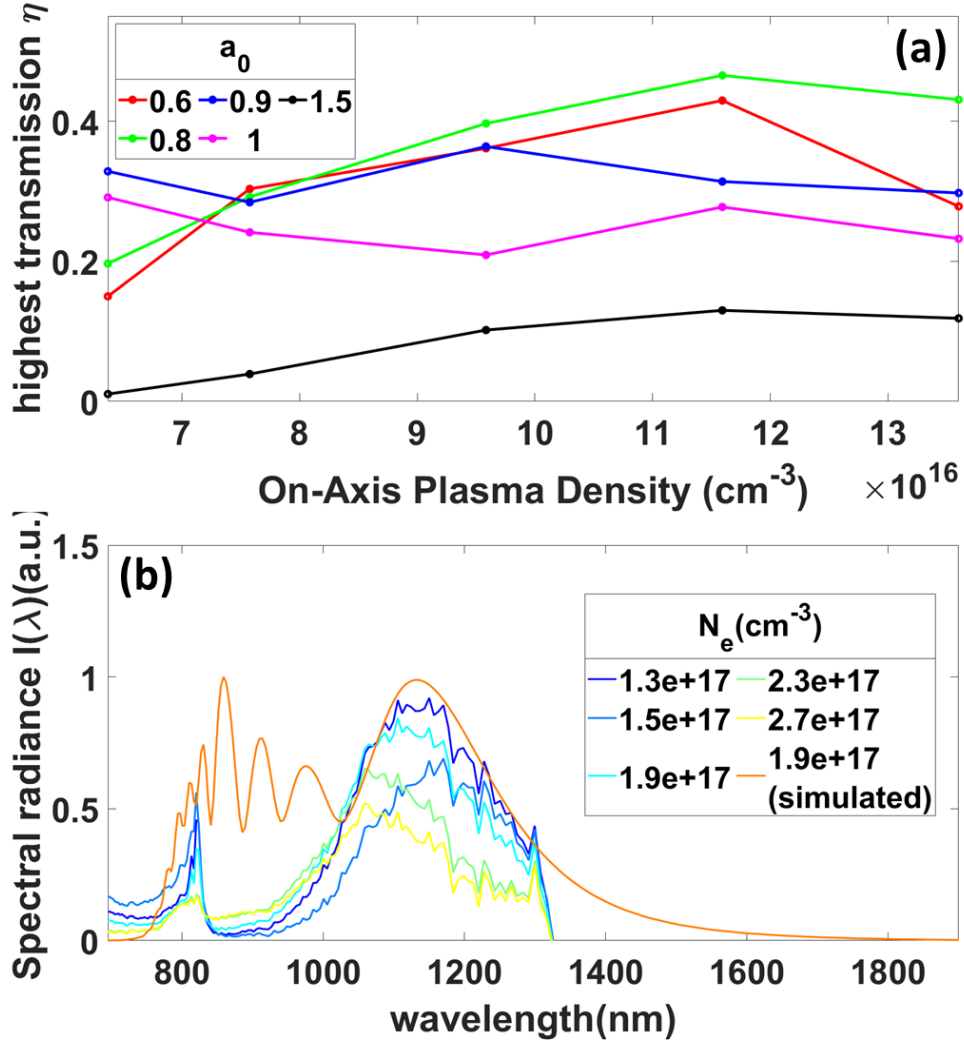


Figure 4.3. a) Energy transmission with different guided beam intensity. (b) Transmitted laser spectrum, measured by Ocean Optics USB2000, after calibration of spectral radiance with an 2800K blackbody reference source (StellarNet SL-4 Cal). The orange curve is the transmitted laser spectra from an FBPIC simulation (see Chapter 4.4 for details).

But besides the error in the measurement introduced by the redshift, the lower apparent transmission efficiency at $a_0 > 0.8$ may also be partially caused by more efficient coupling of laser energy into the plasma waves. The experiments detailed in Chapter 3 and [108], show increased energy transmission as a function of guided beam energy, albeit $a_0 \ll 1$. The efficiency of self-waveguiding might be expected to continue to increase with guided beam energy, as the cost of waveguide formation

is fixed. The efficient coupling of laser energy into plasma waves is ultimately desirable, since the plasma waves accelerate the electron beam, but the more energy that goes into plasma waves, the less that is transmitted by the waveguide. The pump depletion length, discussed briefly in Chapter 1.5, is the characteristic scale length over which the initial pulse energy is deposited in the plasma wave. In the range of intensities in this experiment, as a_0 increases, the coupling efficiency of laser energy into plasma waves increases and the pump depletion length decreases [53], reducing the transmission efficiency. In this experiment, we have no way to separate energy losses from the waveguide and coupling into plasma waves. Future experiments which accurately measure the transmission efficiency and the full laser spectrum, including the IR, can use the redshift to distinguish between energy losses from the self-waveguiding process and from coupling into plasma waves [111].

The guided mode was imaged by a two-lens imaging system and found to be constant within 10% for all but the lowest guided beam intensity and for all on-axis plasma densities (Fig. 4.4). This is to be expected based on estimates of the guided mode size using the measured transverse plasma density profiles in Fig. 3.2 and the Marcuse equation [112], $w/a \approx 0.65 + 1.619/V^{3/2}$ where a is the channel radius, w is the mode size and $V = (2\pi a/\lambda)(n_{core}^2 - n_{cladding}^2)^{1/2}$. Since $V > 10$ for all channels, the guided mode size only depends on the channel radius, which is roughly constant for all guided beam energies and initial plasma densities. The mode sizes here are not consistent with the Marcuse equation, which is likely caused by inaccurately imaging the end of the waveguide, so that the imaged mode has already exited the guide and begun to diffract. The larger mode size for $a_0 = 0.9$ is likely

due to beam being too low intensity at the end of the guide to self-waveguide, causing it to begin to diffract prior to the image plane of the end mode imaging system. The independence of the mode size and on-axis plasma density allows it to be chosen independently of the matched mode size to optimize the electron acceleration.

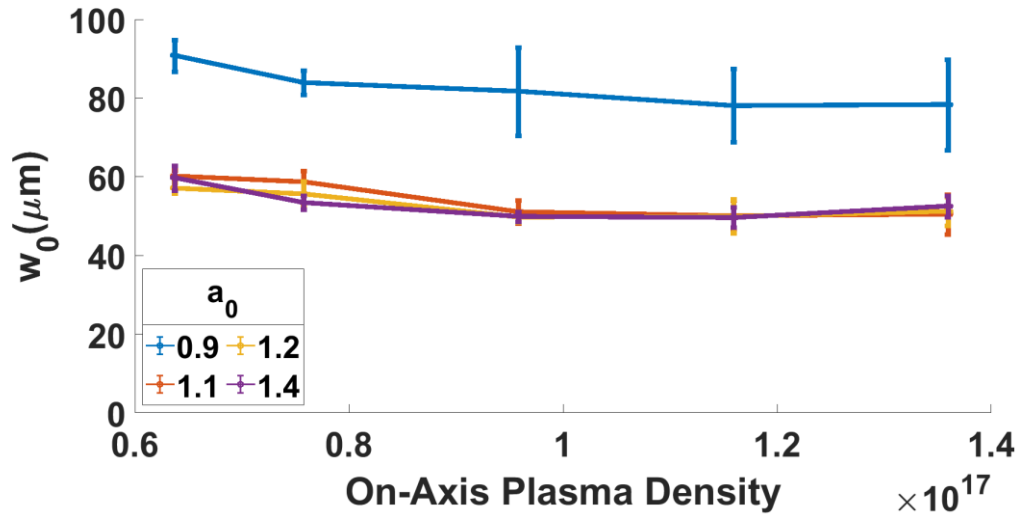


Figure 4.4. Gaussian beam waist of transmitted beam modes after propagation through 20 cm plasma waveguide. The waist corresponds to a fitting the transmitted beam profile to an ellipse and plotting the minor axis. The input beam waist was $35 \mu\text{m}$ (profile shown in Figure 4.3 (b4)). The on-axis plasma density is estimated by scaling of 2-color plasma interferometry (Fig 3.2).

4.4 Acceleration Results

While the redshifted optical spectra indicates that the waveguiding pulse is driving a plasma wave, no electrons were accelerated with only the second gas jet section, since the density and laser intensity was kept low enough to suppress electron self-injection (Chapter 1.4). This was done to restrict injection to the first jet, preventing the broadening of the electron energy spectrum due to continuous injection. To achieve injection in the first section with $a_0 = 1.5$, the first jet section

was doped with 5% argon, to serve as an ionization injection source, as discussed briefly in Chapter 1.4 and demonstrated in [50]. The $a_0 = 1.5$ pulse has a peak intensity of $I = 4.91 \times 10^{18} \text{ W/cm}^2$, well above the barrier suppression ionization threshold of Ar^{10+} , $I_{BSI} = 2.7 \times 10^{18} \text{ W/cm}^2$ [23] (Eq. 1.2). Because of this, the first 11 electrons of Argon and any remaining neutral Hydrogen atoms are fully ionized by the early part of the pulse and acquire the similar momenta as the rest of the background electrons which form the plasma wave. For the remaining ions of the L-shell, Ar^{11-15+} , $I_{BSI} = 4.0 - 23 \times 10^{18} \text{ W/cm}^2$, so most of the ionization of these ions will occur near the peak of the laser pulse, the optimal phase for injection into the wake [113,114]. However, while electrons “born” at the peak of the pulse are most likely to be trapped, ionization at the peak of the pulse is not by itself sufficient to trap electrons. Trapping is also affected by the magnitude of the wake potential, which is itself a function of the plasma density and laser intensity (see Chap 1.3) and by the laser pulse width. If the width of the laser pulse is much shorter than the plasma wavelength, as is the case in our experiment, where $\lambda_p \approx 110 \mu\text{m} > c\tau = 13.5 \mu\text{m}$, then the peak of the pulse occurs farther away from the peak of the wakefield and threshold wake potential required for trapping is increased [114].

The results of the electron acceleration experiment are shown in Figure 4.5. The profile of the electron beam is measured by the fluorescence of a Lanex screen, placed just outside the accelerating chamber after an 850- μm thick aluminum flange. The electrons accelerated in the first section alone are shown in Figure 4.5(b), with a divergence of $\sim 20 \text{ mrad}$ and 90 fC of charge. The energy is too low to be measured

by the spectrometer, but based on the PIC simulations, shown in figure 4.6(a), the energy ~ 100 MeV, with energy spread $\Delta E/E \sim 0.1$.

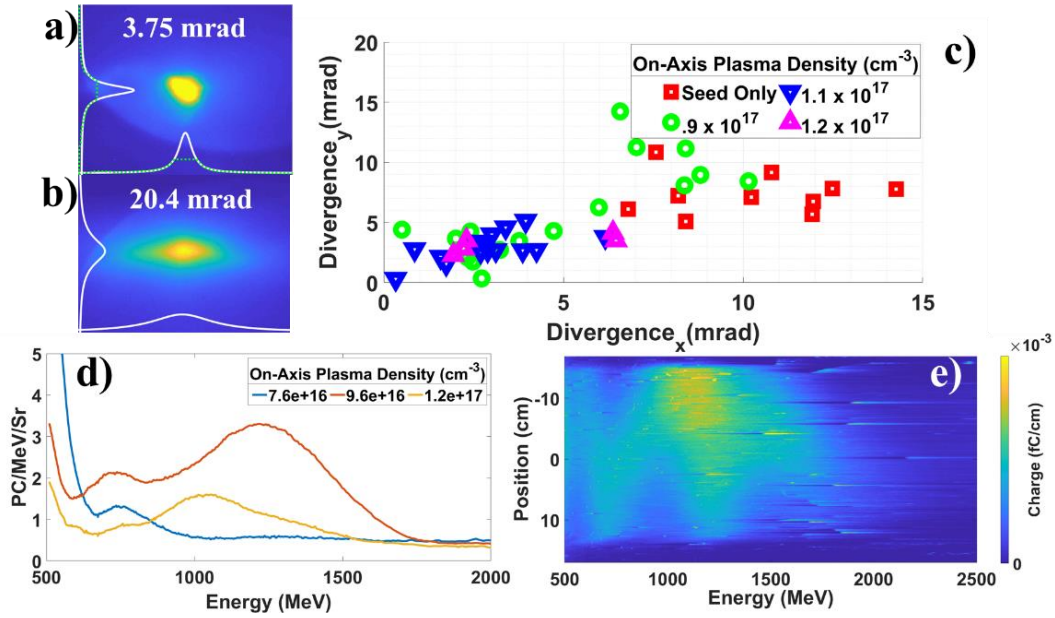


Figure 4.5. (a) shows the electron beam profiles measured by imaging the fluorescence from a LANEX scintillating screen placed in the electron beam 2 meters from the end of the gas jet. The highest energy shot is shown in (a) and the (b) shows the seed beam generated by the injection section alone. The horizontal divergence is labeled in each image. (c) shows the divergence of the electron beam for varying conditions, measured using a Lorentzian fit of the size of electron beam profiles generated by the same LANEX as a,b. (d) Electron spectra from a range of on-axis plasma densities, measured by a permanent magnet electron spectrometer. The spectra shown is the deconvolution of the measured spectrum with the 8 mm slit. (e) Spectrum from the highest energy shot, the red line in (d). For every shot, the J_0 pulse was 1 J, 70 fs and the waveguiding pulse was 8.5 J, 45 fs.

Figure 4.5(a) shows an electron beam profile measured under the same laser conditions, but with the full jet. The fluorescence saturates the camera, but as is shown in the lineout of the beam profile, the saturated image is fitted to a Lorentzian in order to extract a beam width and a divergence. The electron beam divergence extracted using the same method is shown for a variety of conditions is shown in

Figure 4.5 (c). The narrowest beams correspond to the greatest overall energy gain, as the increased longitudinal momentum decreases the beam spread.

After passing through the beam profile Lanex, the electron beam passes through an 8 mm slit and a 30 cm, 1 T magnetic field, permanent magnet spectrometer. The dispersed beam then fluoresces another Lanex screen and is imaged by a camera. Three of the spectra for different initial gas densities are shown in Fig. 4.5(d). The highest energy beams, which were also the lowest divergence, were obtained with $N_0 = 9.6 \times 10^{16} \text{ cm}^{-3}$ on-axis electron density. Integrating the spectrum gives a charge of 4 fC over 1 GeV with the 1.3 GeV peak having approximately 50% energy spread $\Delta E/E$.

In order to understand the acceleration process, particle-in-cell simulations were performed using the PIC code FBPIC [33]. PIC codes simulate the self-consistent interaction of “macroparticles” and field quantities, both defined on a discrete grid and evolved in discrete time-steps. Macroparticles represent many electrons or ions, and have defined positions, momentums, and weights on the grid. During each time-step of the simulation, the fields are calculated then interpolated onto the grid, then the macroparticles are pushed based on the forces from the fields and finally the charge and current are calculated in order to determine the fields for the next step. In FBPIC, the grid is quasi-cylindrical, and the fields are decomposed into azimuthal modes m , [115]

$$F(r, z, \theta) = \text{Re} \sum_{m=-N_m}^{m=N_m} \hat{F}_m(r, z) e^{im\theta} \quad (4.1)$$

This enables fully 3D physics to be modelled at a much lower computation cost than a traditional 3D simulation. Here, the simulation results were found to converge for 3 azimuthal modes, and the total grid size was ($z=6000, r=650, m=3$), more than 200 times less than would have been required for a comparable fully 3D Cartesian simulation. The simulation results shown here were performed with the same parameters as for the best acceleration, $a_0(vacuum) = 1.5$ and $N_{e(on-axis\ average)} = 0.9e17\ cm^{-3}$.

The longitudinal gas density profile of the simulation is the same as in Fig 4.1(b). The transverse profile of the Hydrogen comes from recent 2-color interferometry measurements, taken using the diffractive axicon and a new longitudinal probing technique capable of measuring lower densities. This new technique is still undergoing refinement, but for now it can provide an estimate of the transverse density profile of the channel and help us gain insight into the electron acceleration results.

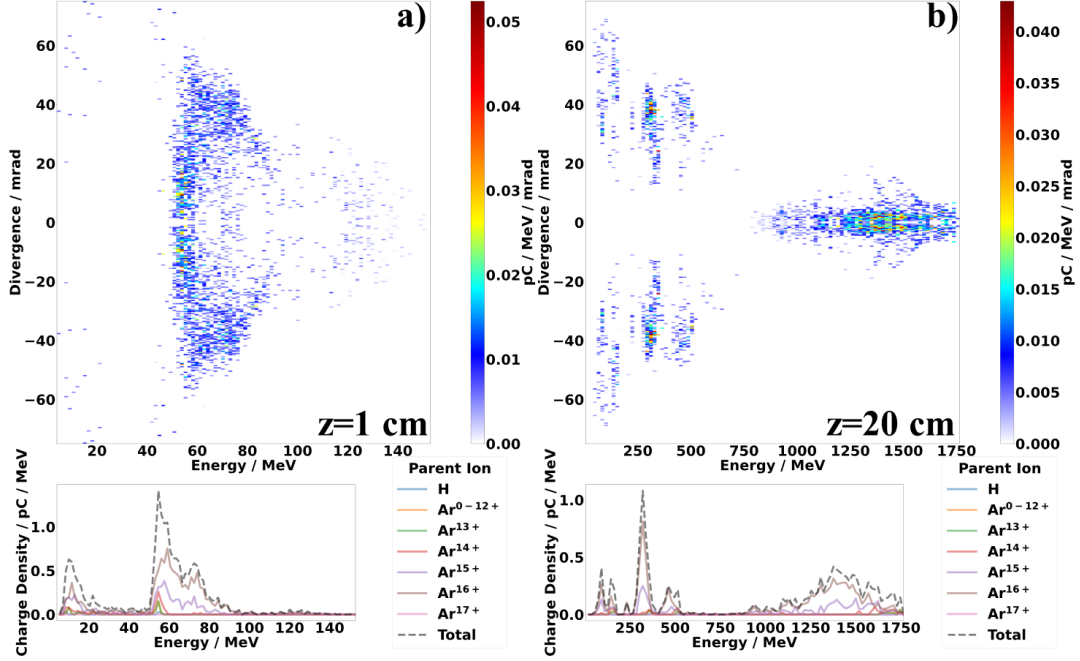


Figure 4.6. FBPIC simulations showing the electron spectra at the end of the injection section (a) and acceleration section (b). The electron spectra integrated over the divergence is shown below, with the electron populations color-coded according to ion parent species. The simulations were performed with the same parameters as the optimal acceleration case for the experiment; $a_0(\text{vacuum}) = 1.5$, $N_{e(\text{on-axis average})} = 0.9e17 \text{ cm}^{-3}$. The longitudinal gas density profile is taken from Fig 4.1(b) and the transverse profile is from recent 2 color interferometry experiments. The simulation was performed in a boosted frame, with $\gamma = 10$ and a preionized plasma.

The simulation results are shown in Figure 4.6. Fig. 4.6 (a) shows the electron beam after the injection section. In the bottom panel, the electrons are color-coded by their parent ion, and only electrons from Ar^{11-15+} are accelerated by the wakefield, with total accelerated charge of 18 pC, average energy of $W = 51 \text{ MeV}$ and divergence of $\sim 35 \text{ mRad}$. Fig. 4.7 (b) shows the same plots after propagation through the whole waveguide. After this section, the accelerated charge has dropped to 14 pC, with most of the lost charge coming from the lower charge Argon ions, Ar^{11-12+} . The electrons from different Argon ions are ionized at slightly different phases with respect to the wake, because of their different ionization potentials. This

can generally lead to broader spectra, but in this case the electrons from the lower charge ions have dephased from the wake over the 19.5 cm propagation. This is likely because of the changing size of the wake due to the longitudinal variations in electron density.

The injection of the electrons that are accelerated through the acceleration section, Ar^{13-15+} , requires some explanation. Rewriting equation 1.2 for the barrier suppression ionization threshold in terms of a_0 gives $a_{0_{BSI}}(Ar^{13,14,15+}) = 1.6, 1.9, 2.1$, above the vacuum a_0 of drive laser. However, the mismatch between the fundamental mode of the waveguide and the laser focal spot leads to mode beating and oscillations in the guided spot size, as is detailed in Chapter 3.3. The beating causes intensity modulations, and the intensity peaks provide regions where ionization of the higher charge Argon ions is possible. Figure 4.7 shows a_0 and the Gaussian beam waist w_0 as a function of propagation distance in the jet. The mode beating also causes changes in the size of the wake, which may contribute to the dephasing of some of the electrons trapped in the injection section.

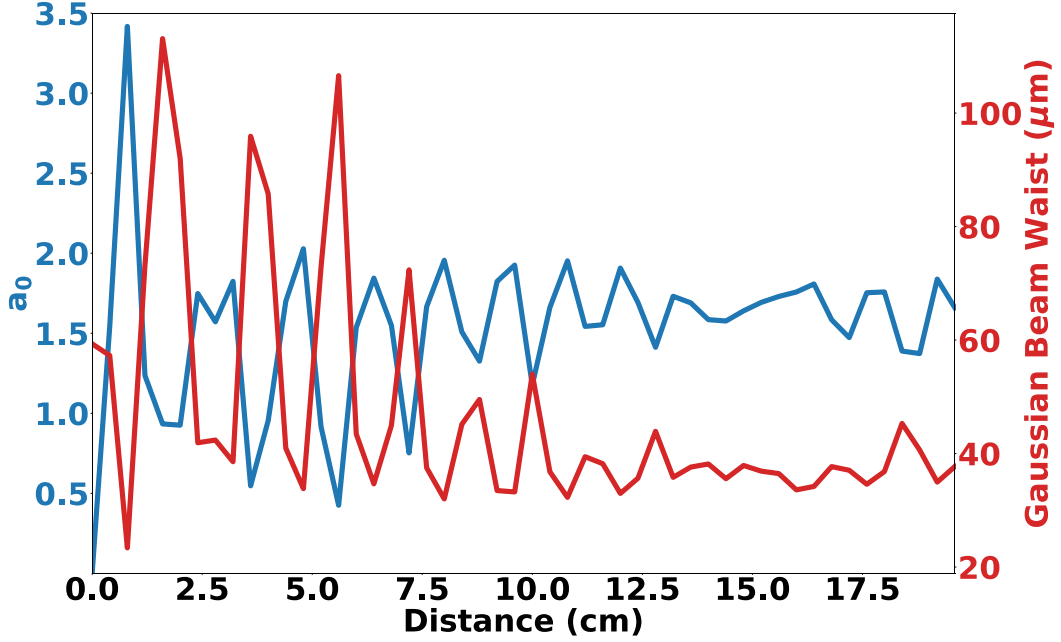


Figure 4.7. Peak a_0 and fitted Gaussian beam waist w_0 as a function of propagation distance for the same FBPIC simulations as Fig. 4.6.

The average energy gain in Fig. 4.6 (b) is $W=1.4$ GeV for the higher energy peak, very close to the experimental result of 1.4 GeV. The energy spread was also similar, $\sim 40\%$ in the simulation vs 50% in the experiment. The major discrepancy between the experiment and the simulation is the accelerated charge, which is 3 orders of magnitude higher in the simulation. The mode beating, as well as changes in the longitudinal size of the wake due to changes in longitudinal density may play a role here, as we have seen from the simulations that trapped some trapped charge is lost in the acceleration section. The experimental transmitted mode is larger than the transmitted mode in the simulation. This could be caused by longitudinal variation in the channel size, which may cause even more loss of charge. Another possibility is that the transverse neutral gas and plasma density profiles need to be more precisely measured to properly simulate the wakefield acceleration.

Pointing jitter between the guided beam and index structuring beam may also be a factor in the low charge. There was a RMS jitter of $4 \mu rad$ in the pointing of the guided beam, which was greater than the jitter in the index structuring beam. This corresponds to a displacement of half the beam waist between the channel and guided beam. Both offsets and pointing mismatches introduce higher order modes into the guide and the beating between these modes and the fundamental mode of the waveguide causes spot size oscillations. The higher order modes have a much high leakage rate and reduced coupling into the accelerating plasma waves [116].

Although the energy gain in the simulations agrees with the experiments it is significantly lower than analytic estimates. Finding the expected energy gain with Eq. 1.8 using the average on axis plasma density as N_{e0} yields 3.3 GeV, close to double the value from the experiment. The mode beating and other spot size oscillations may again play a role here, as the higher order waveguide modes are likely less efficient at driving the wake. Further simulations are necessary to understand all these effects for future experiments. Measurement of the plasma wave amplitude using the redshift of the optical spectra could also help us understand the acceleration process and provide another point of comparison for the simulations.

Here we have shown the results of the first all-optical guided LWFA with electron energy gains of over 1 GeV. We have extended the self-waveguiding technique to $a_0 > 1$ and 20 cm length. We have also demonstrated a new sectioned gas jet which was used to localize ionization injection. Simulations of the process have offered insight into the injection and acceleration process. Future experiments will be able to utilize the flexibility and adaptability of both the gas jet and the self-

waveguiding technique to further optimize and design bigger and better laser wakefield accelerators.

Chapter 5: Future Work

5.1 Summary

The long-term goal of the research in this dissertation is the development of a laser wakefield accelerator with sufficiently high energy and beam quality for high energy physics applications. Much attention has been focused on the potential for achieving that goal through combining multiple quasi-linear LWFA modules of ~ 10 GeV energy gain per stage [17,52,117]. A key component required for these modules is a meter-scale plasma waveguide to extend the length of the accelerator beyond the Rayleigh range of the drive laser. The plasma waveguide also must have low enough on-axis density ($N_{e0} \sim 10^{16-17} \text{ cm}^{-3}$) to prevent dephasing of the electron beam and the wakefield. Chapters 2 and 3 present two new techniques for creating these plasma waveguides: the two Bessel and self-waveguiding methods.

In the two Bessel method, a J_0 Bessel beam formed by an axicon ionizes a plasma column. The plasma column then expands hydrodynamically, leaving the low on-axis plasma density required for wakefield acceleration. A second pulse passes through a spiral phase plate and is focused by the axicon into a J_{16} or J_8 Bessel beam. The higher order Bessel beam then ionizes the ring of neutral gas which had been pushed aside by the expanding plasma, creating the plasma waveguide. This technique was demonstrated over 5 cm for a guided pulse of $I = 6 \times 10^{16} \text{ W/cm}^2$ and for 30 cm with a lower power guided pulse.

In the self-waveguiding method, the neutrals on the edge of the expanded plasma column are instead ionized by the front of the guided pulse itself. Proof-of-

principle experiments conducted at University of Maryland and presented in Chapter 3 demonstrated guiding of $I = 3 \times 10^{17} W/cm^2$ over 10 cm.

Chapter 4 presented the results of a preliminary effort to use the self-waveguiding technique in a full LWFA. The experimental campaign at Colorado State University's ALEPH laser extended the waveguide to 20 cm and guided pulses of $I = 1.5 \times 10^{18} W/cm^2$. This experiment also added a short, Argon doped section to the gas jet target and injected electrons in the waveguide, achieving acceleration of over 1.5 GeV.

5.2 Future Work

Both the self-waveguiding and two-Bessel methods have shown great guiding results at lower guided beam intensities than are required for electron acceleration. Chapter 3.4 compared the two methods in this regime, but a clear next step for the two-Bessel method is guiding a beam with $a_0 \sim 1$, a high enough intensity for LWFA. That will enable a more direct comparison of the two methods for the regime in which they will be applied. It is possible that the higher intensity guided beam will ionize any neutral gas near the waveguide not ionized by the higher order Bessel beam, forming a similar waveguide at a similar energy cost to the self-waveguiding method.

Either method could be improved by more accurate and higher resolution measurements of the plasma and neutral gas expansion, particularly at lower initial gas densities. Such measurements will be necessary to best tune the waveguides to match the fundamental mode of the waveguide to the focal spot of a high-power guided beam for wakefield acceleration.

The preliminary electron acceleration experiments in Chapter 4 demonstrate the feasibility of GeV scale electron acceleration using the self-waveguiding method, but there are a few ways in which the experiment could be changed to produce higher energy and lower emittance electron beams. There are also still questions about the results which further experiments with a similar setup could help answer.

Beginning with the self-waveguiding process, transmission measurements sensitive to the infrared are necessary to characterize the guiding efficiency for different initial gas densities and laser intensities. Analysis of the redshifting of the transmitted beam can also be used to separate any losses from the self-waveguiding process compared to laser energy invested into plasma waves [111].

Better knowledge of self-waveguiding at high intensities can be used to optimize the electron acceleration. The energy gain can also likely be improved by extending the length of the acceleration; the 20 cm waveguide used in Chapter 4 is far shorter than the 70 cm dephasing length for the optimal acceleration conditions found in the experiment [14]. A higher intensity waveguiding pulse would also increase the energy gain, and may dramatically increase the accelerated charge, as is discussed in Chapter 4.4.

In terms of emittance and energy spread, FBPIC simulations were performed which show that reducing the longitudinal variation of the waveguide can reduce the divergence of the electron beam by an order of magnitude. The energy spread could be decreased by tuning the dopant concentration and injection section density, or by using an alternative injection method such as shock injection.

Moving beyond the basic structure of the Chapter 4 experiment, the inherent tunability and flexibility of self-waveguiding and two-Bessel waveguides opens a large new parameter space for LWFA. Some possible future experiments which have been explored in simulations and are enabled experimentally by these waveguides include using smaller, higher density waveguides to access the bubble regime in channels [118], lower density, nearly hollow channels for positron acceleration [119] and longitudinal density tailoring to preserve beam emittance [120]. This guiding techniques discussed in this thesis are a promising new alternative to the capillaries or self-guiding used in most LWFA experiments in the last twenty years, and are likely to become a key part of any future high energy laser-driven wakefield accelerators.

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