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Channeling of intense laser beams in underdense plasmas

M. D. Feit and J. C. Garrison
Lawrence Livermore National Laboratory, University of California, Livermore, California 94550

A. M. Rubenchik
Department of Applied Science, University of California, Davis, Livermore, California 94550

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A hydrodynamic simulation is used to show that intense laser pulses propagating in underdense plasmas create stable, long-lived, and completely evacuated channels. At low intensities, $I=10^{17}$ W/cm^2, self focusing seriously distorts the temporal envelope of the pulse, but channeling still occurs. At high intensities, $I=10^{19}$ W/cm^2, channeling can proceed over many diffraction lengths with significant distortion restricted to the leading edge of the pulse. [S1063-651X(97)50709-4]

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The recent development of petawatt lasers has made it possible to generate high-intensity long-duration pulses that can create a completely evacuated channel in an underdense plasma. A stable channel, which could serve as an optical waveguide, would be a significant contribution to the success of the fast ignitor concept [1] for inertial confinement fusion, since this scheme requires an intense laser beam to propagate a long distance without significant diffraction. In this paper we present the first hydrodynamic simulation of laser-induced cavitation in underdense plasmas which includes all of the following physical effects: relativistic electron dynamics, wave propagation, and ion motion. Our model provides a self-consistent description of ion and electron expulsion by the laser ponderomotive force, including the formation of strong shocks, together with the propagative effects of self-focusing, caused by the relativistic increase in electron mass and by the expulsion of the charged particles. For the laser pulse parameters needed by the fast ignitor, $I_L \geqslant 10^{17}$ W/cm^2 and $\tau_L \approx 100$ ps, both focusing mechanisms are important; therefore, it is essential to study the possible destruction of the channel by these effects. Stable channel formation [2] and destructive filamentation [3] have both been reported, and an adequate simulation would be useful for understanding these experiments. Electron cavitation, i.e., the expulsion of electrons only from the region of high field, has been studied previously for short pulses (duration $\tau_L \approx 100$ fs) using the assumption that ion motion could be completely neglected [2,4–7]. The effects of ion motion for short pulses ($\tau_L \approx 30$ fs) and short propagation lengths ($z \approx 40$ μm) have been studied with a 3D particle-in-cell simulation [8]. Previous hydrodynamic calculations, including ion motion, either used nonrelativistic dynamics for the electrons and ray tracing for the propagation of the laser pulse [3,9,10], or assumed a prepulse which created the channel in which propagation took place [11,12]. These approaches are not suitable for the fast ignitor that requires high intensities and small spot sizes. The application of averaging techniques to the combined hydrodynamic and wave propagation equations makes the treatment of much longer pulses ($\tau_L \approx 100$ ps) and propagation lengths ($z \approx 90$ μm) computationally tractable. With this model, we show that a single pulse can induce complete channeling, i.e., the expulsion of both electrons and ions from the field region to form a so-called ''vacuum tube'' [12] which persists long after the laser pulse has passed.

Long-pulse propagation in an underdense plasma is characterized by several time scales: $\tau_L > \tau_i \gg \tau_e \gg 1/\omega$, where τ_L is the pulse duration, τ_i and τ_e are respectively the ion and electron response times, and ω is the laser frequency. The corresponding spatial scales in the plasma are much larger than the laser wavelength. We model this situation by the propagation of an intense pulse through a cold collisionless plasma having a single species of ions with mass m_i and charge $Z_i e$. For the contemplated intensities, the electrons acquire relativistic quiver velocities within a few optical periods; therefore, they will be modeled as a relativistic cold fluid, i.e., all electronic thermal effects will be neglected. The massive ions will form a cold nonrelativistic fluid. The relation $\tau_i \gg \tau_e$ shows that the (nonquiver) electron motion is adiabatic with respect to the ions; therefore, the multiple-scales analysis of the electron fluid equations for short pulses [13] can be generalized to long pulses by simply allowing for a nonzero ion velocity. This does not change the form of the electron fluid equations, so the short pulse analysis of the electron motion is directly applicable to the long pulse case.

In describing the short pulse results for electrons, it is convenient to use the dimensionless propagation variables $\tau' = \omega_P t$ and $\zeta' = k_P(z-ct)$, where $\omega_P = \sqrt{4\pi N_{e0} e^2/m}$ is the plasma frequency, $N_{e0}(N_{i0})$ is the undisturbed electron (ion) density, and $k_P = \omega_P/c$. We have shown [13] that the electron fluid approaches a quasistatic state, i.e., independent of τ' [14], after an interval of a few plasma periods. Since $\tau_i \gg 1/\omega_P$, this conclusion still holds in the long pulse case. Furthermore the condition $c\tau_L \gg 1/k_P$ implies that the ζ' dependence of the electron fluid variables will be weak compared to their dependence on the transverse coordinates. Thus neglecting both the τ' and ζ' dependence of the electron fluid variables is a good approximation for the long pulse problem. This leads to a model, previously used in treatments of the short pulse problem [2,15], in which the balance between electrostatic and ponderomotive effects is expressed by $\phi = \gamma - 1$, where $\phi = e\Phi/mc^2$ is the normalized electrostatic potential, the electron γ-factor is given by $\gamma=\sqrt{1+|a_L|^2/2}$, and $a_L=eA_L/mc^2$ is the normalized, slowly varying amplitude of the laser field. Combining these relations with the Poisson equation yields the electron density as

$$n_e=\max\left[\frac{n_i+\nabla_\perp^2\gamma}{\gamma},0\right], \tag{1}$$

where $n_e=N_e/\gamma N_{e0}$ is the normalized (proper) electron density, $n_i=N_i/N_{i0}$ is the normalized ion density, $\nabla_\perp$ is the transverse gradient, and the max function enforces the physical restriction of non-negative density [2,15].

The quasistatic description of the electrons must be supplemented by the fluid equations for the ions. Since the ion response time can be comparable to the pulse duration, the τ' and ζ' dependence of the ion variables cannot be ignored, but other simplifications can be made. Reverting to laboratory coordinates, we note that the longitudinal component of the ponderomotive force is $F_z=-mc^2\partial_z\gamma=-mc^2\partial_z|a_L|^2/(4\gamma)$. In the limit of long pulse duration the laser pulse is approximately constant in z, so $F_z\approx 0$. Thus a reasonable approximation is to neglect the z component of the ion velocity and to average the ion equations in z over a length small compared to the pulse length, but large compared to the wavelengths associated with ionic excitations. This is equivalent to neglecting all z-derivatives in the ion fluid equations. In formulating this model we use ''plasma units'' in which $m=c=\omega_P=1$, i.e., all lengths are measured in units of $1/k_P$. Then using the dimensionless coordinates $\tau=t-z$ and $\zeta=z$, the simplified equations are

$$[\nabla_\perp^2+2ik\partial_\zeta]\mathbf{a}_L=\left(n_e+\frac{1}{M}n_i\right)\mathbf{a}_L, \tag{2}$$

$$\partial_\tau\mathbf{u}+\frac{1}{2}\nabla_\perp\mathbf{u}^2=-\frac{1}{M}\nabla_\perp\gamma, \tag{3}$$

$$\partial_\tau n_i+\nabla_\perp\cdot(n_i\mathbf{u})=0, \tag{4}$$

where k is the laser wave number in plasma units, $M=m_i/(mZ_i)$, and $\mathbf{u}$ is the ion velocity in units of c. The neglect of thermal effects can be justified by the observation that the preformed plasma is expected to have a temperature of at most a few keV. Thus the thermal pressure will be negligible compared to the ponderomotive pressure, and the laser energy will be deposited in the plasma in the form of fluid motion rather than heating.

Before discussing the simulation results, it is instructive to make some rough estimates. In conventional units the transverse ion acceleration is $\alpha_\perp=-mc^2\nabla_\perp|a_L|^2/(4M\gamma)$, so $\alpha_\perp\approx mc^2|a_L|^2/(4M\gamma w)$, where w is the beam radius. Consequently the transverse velocity after a time τ is approximately $v(\tau)=mc^2\tau|a_L|^2/(4M\gamma w)$. The time required for expulsion of ions ($\tau_{\rm cav}$) is estimated by assuming constant acceleration from rest, so that $v(\tau_{\rm cav})\tau_{\rm cav}=2w$ and $c\tau_{\rm cav}=2\sqrt{M\gamma/m}w/|a_L|$. The velocity at this time is $v_{\rm cav}=c|a_L|/(2\sqrt{M\gamma/m})$. For the beam radius $w=7$ μm, mass ratio $M/m=7200$, and intensity range 10^{19} W/cm$^2\geqslant I\geqslant 10^{17}$ W/cm^2, one finds 4 ps $\leqslant\tau_{\rm cav}\leqslant 40$ ps and $0.01\geqslant v_{\rm cav}/c\geqslant 0.001$. The cavity will continue to expand for some time after the pulse has passed, and in this limit the thermal pressure will become important. We estimate the final cavity radius by equating the energy deposited by the pulse to the PdV work done in creating a cavity of radius R in a fluid with temperature $T=1$ keV (a typical temperature for the preformed plasma), with the result $R/w=\sqrt{mc^2/(8k_BT)}|a_L|$. Then $|a_L|=O(1)$yields $R\gg w$, i.e., the cavity radius will be larger than the spot size of the beam. After the pulse has passed, the cavity will collapse in the time $T_{\rm clps}=R/v_{\rm th}$, where $v_{\rm th}=\sqrt{k_BT/M}$. With the nominal temperature $T=1$ keV and the values considered below the collapse time is $T_{\rm clps}\approx 100$ ps. By contrast, the electrons-only cavity, which requires a charge separation between electrons and ions, will collapse as the pulse is passing. The energy stored in the electrons is thereby returned to the field which experiences no loss during propagation. In the case of complete cavitation, the rate of energy loss from the pulse is estimated by equating the difference in intensity at two

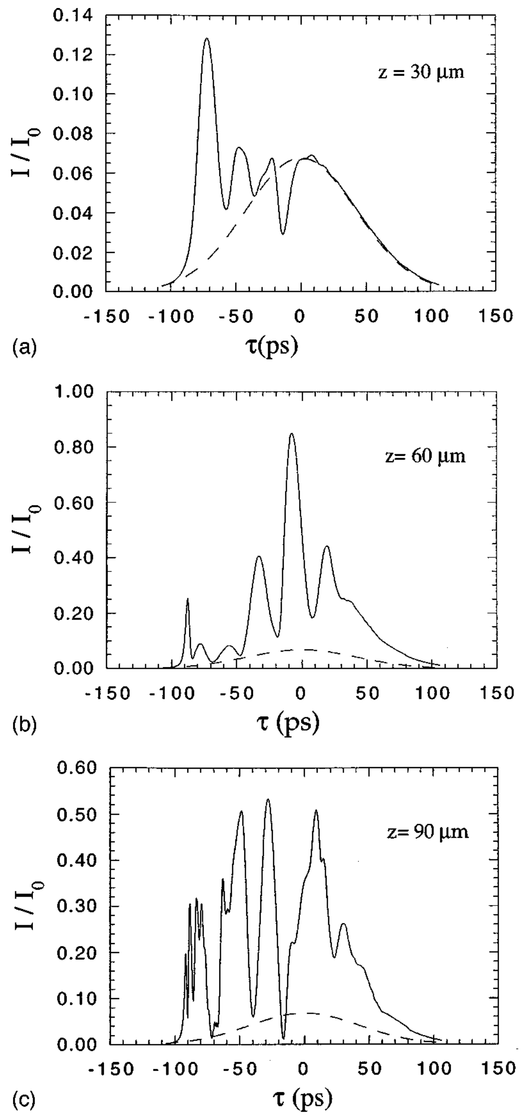

FIG. 1. Normalized laser intensity on axis vs $\tau=t-z/c$, at $z=$(a) 30, (b) 60, and (c) 90 μm, for an incident pulse with $I=10^{17}$ W/cm^2 and a temporal profile shown by the dashed curve. The normalization intensity $I_0=1.48\times 10^{18}$ W/cm^2 corresponds to $|a_L|=1$.

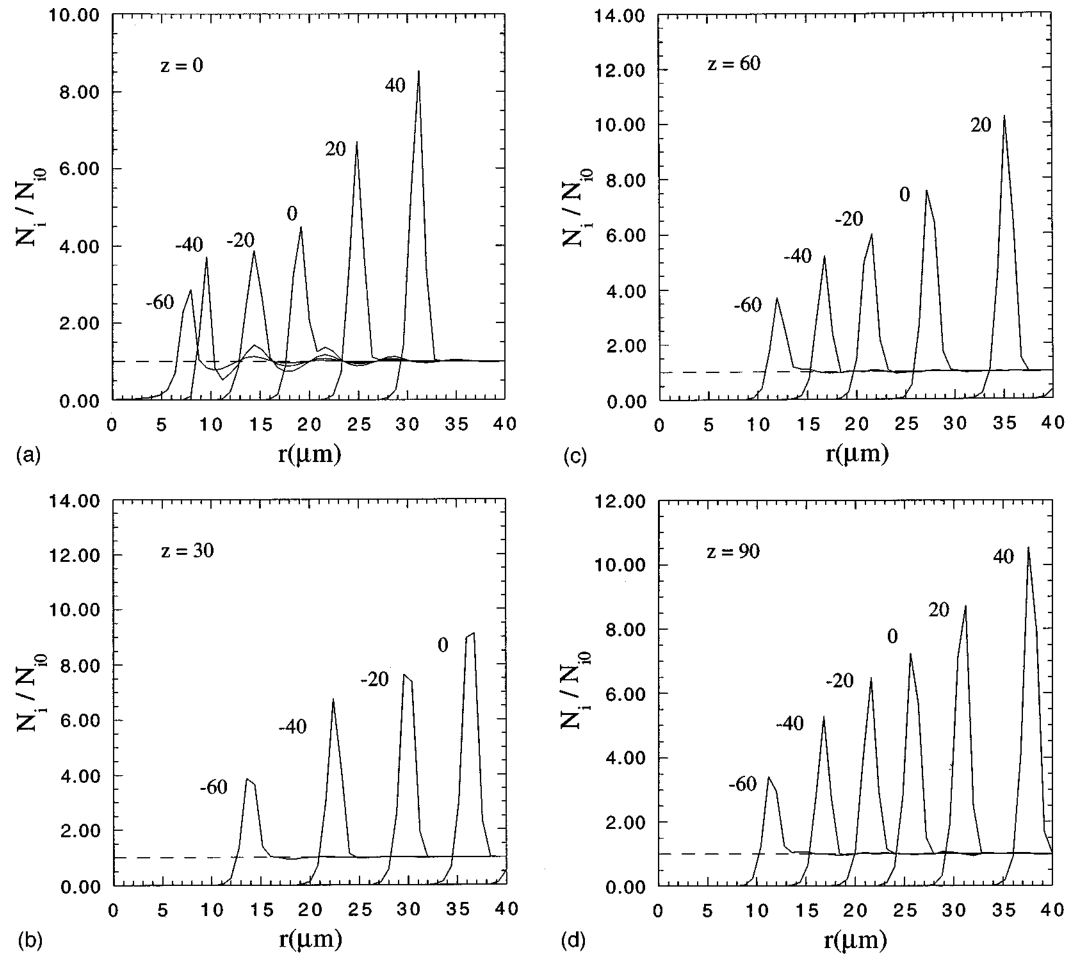

FIG. 2. Snapshots of the normalized ion density $n_i = N_{i/N_{i0}}$ vs r (μm) at several values of τ for $z=$(a) 0, (b) 30, (c) 60, and (d) 90 μm, for the same case as Fig. 1. The horizontal dashed line represents the initial density N_{i0}, and the number next to each curve gives the snapshot time in ps.

planes, separated by Δz, to the rate of energy deposition in the cylindrical volume defined by the planes; i.e.,

$$dI_L/dz = -n_i(Mv_{\text{cav}}^2/2)/\tau_{\text{esc}}, \tag{5}$$

where the escape time is $\tau_{\text{esc}} = w/v_{\text{cav}}$. In terms of $|a_L|$ this equation is

$$|a_L|^{-2} d|a_L|^2/dz = -(1/L_0)(|a_L|/\gamma^{3/2}), \tag{6}$$

where $L_0 = 8(n_{\text{cr}}/n_i)\sqrt{M/m}w$, and n_{cr} is the critical density for the laser frequency. Thus the absorption length is $L_{\text{abs}} = \gamma^{3/2}L_0/|a_L|$. We compare L_{abs} to L_R, the characteristic diffraction length for transverse modulation on the scale of the plasma wavelength, which is defined by $L_R = kk_P^{-2} = (n_{\text{cr}}/n)k_P^{-1}$. The result is

$$L_{\text{abs}}/L_R = 8\sqrt{M/m}k_P w\gamma^{3/2}/|a_L| \geqslant 9.12\sqrt{M/m}k_P w \gg 1, \tag{7}$$

i.e., the pulse can propagate for many diffraction lengths before suffering significant loss.

The propagation equation (2) is solved by a spectral method for paraxial propagation [16]. The momentum-balance equation (3) is approximated by a finite-difference scheme using upwind differencing, and the continuity equation (4) is treated by integrating it over concentric spherical shells surrounding the radial grid points, with fluxes defined by the same upwind scheme. In the results presented below we impose cylindrical symmetry and use the incident laser field $a_L = a_{L0}\exp[-r^2/(2w^2)]\exp[-\tau^2/(2T_P^2)]$. The laser wavelength, spot size, and pulse duration are respectively $\lambda = 1.0$ μm, $w = 7$ μm, and $T_p = 60$ psec in all cases. The initial electron density is $n_e = 0.1$ n_{cr}, which satisfies the necessary condition $n_e \ll n_{\text{cr}}$.

For the lowest intensity, $I = 10^{17}$ W/cm^2, strong self-focusing effects can be seen in Fig. 1, which shows the temporal profile of the laser intensity on axis for three values of the propagation depth. The temporal modulation, which is similar to results obtained in calculations for Kerr media [17], is a consequence of the variation in self-focusing lengths during the pulse. There are also important differences from the case of Kerr media; in particular the formation of the channel prevents truly singular self focusing. The severe pulse distortion shown in Fig. 1 does not prevent the formation of a cavity. This is seen in Fig. 2, which shows the ion density as a function of radius, r, at various times, $\tau = t - z/c$, for the same propagation depths. The time required for the formation of the cavity and the cavity radius agree with the estimates, and the shock front forming the wall of the cavity at each time is clearly evident. The plots in Fig. 2 show that the increased intensity due to self-focusing

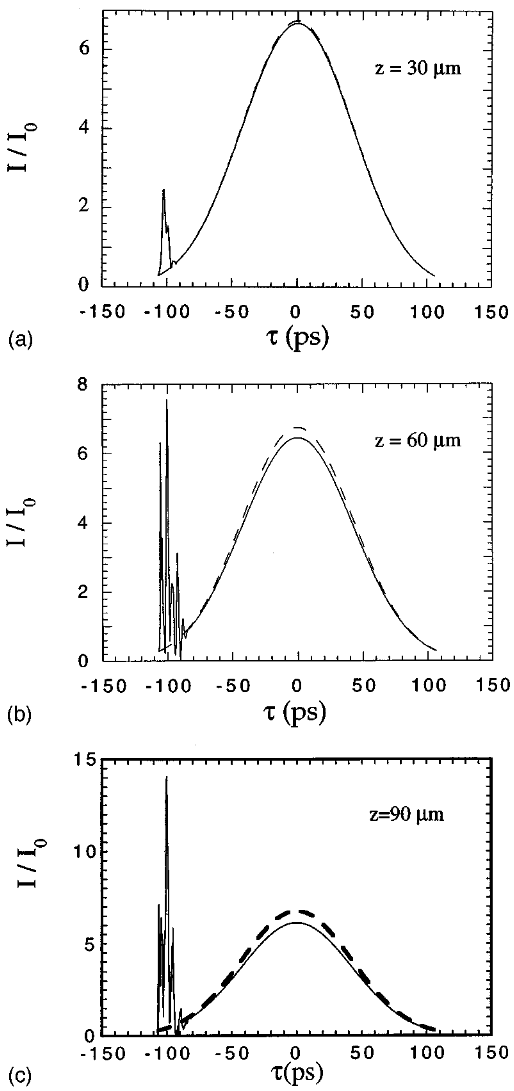

FIG. 3. Normalized laser intensity on axis vs $\tau = t - z/c$, at $z = $(a) 30, (b) 60, and (c) 90 μ for an incident pulse with $I = 10^{19}$ W/cm^2. Other parameters as in Fig. 1.

causes more rapid cavity formation as the pulse propagates, but the cavity radius at a given time τ may be decreased somewhat during propagation, as seen by comparison of the plots for $z = 30$ and 60 μm.

At the highest intensity, $I = 10^{19}$ W/cm^2, the destructive effects of self-focusing are largely suppressed, as seen in Fig. 3. The temporal modulation of the pulse increases with propagation depth, but it is now confined to the leading edge where the intensity is low. In this case self focusing is self-limiting, since the increase in the intensity at the leading edge drives rapid cavity formation so that the following part of the pulse effectively propagates in a vacuum and experiences no self focusing. This is an improvement over the low intensity behavior, but the pulse shape will eventually be degraded at any intensity. The fraction of the pulse experiencing modulation increases by roughly 2% for each 10 μm of propagation length, so even the high intensity pulse will be severely deformed after propagation of several hundred micrometers. Cavity formation is similar to the low intensity case, but occurs more quickly, in line with the estimates.

These simulations show that intense laser pulses propagating in an underdense plasma will produce completely evacuated and long-lived channels. The longitudinal profile of the pulse is distorted by self focusing, but at higher intensities the rapid formation of the cavity tends to suppress this effect. Our results support a fast ignitor concept in which an intense pulse of a few ps duration creates an evacuated channel which will persist for $T_{\mathrm{clps}} \approx 100$ ps. The heating pulse can then propagate in the channel without loss or distortion. In a subsequent publication we will present a more detailed account including applications to non-Gaussian beam shapes.

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