Electromagnetic Energy Penetration in the Self-Induced Transparency Regime of Relativistic Laser-Plasma Interactions

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Two qualitatively different scenarios for the penetration of relativistically intense laser radiation into an overdense plasma, accessible by self-induced transparency, are presented. In the first one, penetration of laser energy occurs by solitonlike structures moving into the plasma. This scenario occurs at plasma densities less than approximately 1.5 times the critical one (depending on ion mass). At higher background densities, laser light penetrates only over a finite length which increases with incident intensity. In this regime the plasma-field structures represent alternating electron (and, on longer time scales, ion) layers separated by about half a wavelength of cavitation with concomitant strong charge separation.

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Recent developments of laser technology have opened possibilities to explore laser-matter interactions in regimes previously not achievable [1]. This has given a strong impulse to theoretical investigations of phenomena occurring in such extreme conditions, when electrons quiver with relativistic velocities and new regimes may appear. In particular, penetration of ultraintense laser radiation into overdense plasmas plays a fundamental role in the development of the fast ignitor fusion concept as well as of x-ray lasers [2,3]. In connection with such a problem, many detailed studies of intense laser-plasma interactions have recently been carried out based on numerical simulations with relativistic particle-in-cell (PIC) codes [3–6], multifluid plasma codes [7], and Vlasov simulations [8], as well as experiments [9].

In this Letter, we focus on a fundamentally new feature of the relativistic laser interaction with overdense plasmas in the so-called self-induced transparency (SIT) regime where the optical properties of the plasma are substantially modified by the relativistic increase of the inertial mass. As was shown in the 1970s, in this regime, relativistic effects enable super-intense electromagnetic radiation to propagate through classically overdense plasmas [10–12]. Up-to-date, the main conclusions about this process are as follows: First, there is a threshold intensity for penetration [11]; for intensities lower than this threshold, an overdense plasma reflects the incident radiation with the formation of a nonlinear skin-layer structure at the plasma-vacuum boundary. Second, for higher intensities, the radiation was found to propagate in the form of nonlinear traveling plane waves [10], or solitary waves with specific properties [13]. On the other hand, an exact analytical study recently done of the stationary stage of the penetration of relativistically strong radiation into overdense plasma, taking into account both the relativistic and striction nonlinearities, has led to the discovery of another scenario that is also possible for incident intensities exceeding the threshold [14]. In fact, if the background plasma density N_o is less than $1.5N_c$ ($N_c = m\omega^2/4\pi e^2$ is the critical density and ω is the laser frequency), the stationary solutions localized in the direction of propagation exist only in the form of nonlinear skin layers. This means that the penetration allowed by the SIT effect, if it happens, must have a dynamical nature. At higher densities $N_o > 1.5N_c$, due to electron cavitation, the penetration occurs over a finite length only. In this case, the subsequent plasma-field structure consists of alternating electron layers, separated by depleted regions with an extension of about the wavelength and acting as a distributed Bragg reflector. We emphasize that the nature of this cavitation has a direct analogy with 2D cavitation in the case of relativistic self-focusing [15] where new multifilament structures can also be created [16]. The goal of this Letter is to provide insight into this problem and to find an answer to what kind of scenarios, of electromagnetic energy penetration/transmission into overdense plasmas, may take place.

Here we employ a relativistic cold fluid model assuming that the particle energies of the regular motion are much larger than the electron and ion temperatures. The fluid model represents a significant simplification over a full kinetic treatment based on particle-in-cell simulations or the Vlasov equation, but retains enough physics to be qualitatively and quantitatively useful. This simplification comes at the price of having discarded most of the kinetic behavior such as fast electron and ion beam production where, however, only a few percent of the incident laser energy is deposited. The next step of simplification is to use an averaging procedure to remove the fast time scale associated with the laser carrier frequency. In this case, the governing set of self-consistent equations for the one-dimensional problem in the Coulomb gauge reads (see, for example, [17])

$$\frac{\partial p_{\parallel e,i}}{\partial t} = -q_{e,i}\frac{\partial \varphi}{\partial x} - \frac{\partial \gamma_{e,i}}{\partial x}, \tag{1}$$

$$\frac{\partial n_{e,i}}{\partial t} + \frac{\partial}{\partial x}\left(n_{e,i}\frac{p_{\parallel e,i}}{\gamma_{e,i}}\right) = 0, \tag{2}$$

$$\frac{\partial^2 \varphi}{\partial x^2} = n_o(n_e - Zn_i), \tag{3}$$

$$\frac{\partial^2 a}{\partial x^2} + \left[1 - n_o\left(\frac{n_e}{\gamma_e} + \frac{Z^2 n_i}{\gamma_i}\right)\right]a - 2i\,\frac{\partial a}{\partial t} = 0. \tag{4}$$

Variables are normalized as follows: $\omega t \to t$, $\omega x/c \to x$, the longitudinal momentum of the electrons and ions $p_{\parallel e,i}/mc \to p_{\parallel e,i}$; the scalar potential $e\varphi/mc^2 \to \varphi$; electron and ion densities $N_{e,i}/N_o = n_{e,i}$; $\gamma_e = (1 + p_{\parallel e}^2 + a^2)^{1/2}$ is the electron Lorentz factor, $\gamma_i = [\delta^2 + (p_{\parallel i}^2 + Z^2a^2)]^{1/2}$; $n_o = N_o/N_c$ is the overcritical parameter; $\delta = M/m$ is the ratio of ion and electron mass; $q_e = -1$ and $q_i = Z$ are the electron and ion charges; and we consider circularly polarized laser radiation with the vector potential taken as $e\mathbf{A}/mc^2 = \mathrm{Re}[a(x,t)(\mathbf{e}_y + i\mathbf{e}_z)\exp(i\omega t)]$.

Equations (1)–(4) have been numerically integrated for the situation of normally incident laser radiation from vacuum ($x < 0$) onto a semi-infinite overdense plasma ($x \geq 0$), the numerical interval consisting of two parts: a short vacuum region ($-x_b \leq x < 0$) to the left of the plasma boundary, and a semi-infinite plasma region to the right. As for the boundary conditions, at infinity the fields must vanish ($a = 0$, $\varphi = 0$), particles are immobile ($p_{\parallel e,i} = 0$), and the electron and ion densities are unperturbed ($n_{e,i} = 1$). In the vacuum region (at $x = -x_b$) the radiation boundary condition for the laser field reads

$$a + i\,\frac{\partial a}{\partial x} = 2a_i(t), \tag{5}$$

where $a_i(t)$ is the amplitude of the incident laser wave, which means that in the vacuum region the total field is the sum of the incident and reflected wave. At the initial time the electrons are in equilibrium with the ions, i.e., $p_{\parallel e,i} = 0$, $n_{e,i} = 1$, $\varphi = 0$. The analysis has been performed for overdense plasmas ($n_o > 1$) and for a quite wide range of incident intensities.

In order to emphasize the key role played by the ponderomotive force in the propagation dynamics, we first focus on the situation with fixed ions ($\delta = \infty$). In order to make sure that the code works properly, we start the simulations with a semi-infinite pulse turning on as a *tanh* function. In this case, for maximum incident intensities lower than the threshold of penetration, after a transient stage, a stationary regime with the formation of a nonlinear skin layer is reached, which is in perfect agreement with the previous analytical solutions. Furthermore, good agreement is found with the calculated threshold for penetration [11]. For intensities above the threshold, the nonlinear skin-layer regime is broken and the interaction leads to penetration of laser energy into the overdense plasma. The analysis of the dynamical process, above the threshold, has revealed two qualitatively different scenarios of laser penetration, depending on the supercritical parameter n_o. The same qualitative behavior of the system occurs over a wide range of incident intensities and thus it does not sensitively depend on the specific values.

If $n_o \leq 1.5$, we have only a dynamical regime where laser radiation slowly penetrates into the overdense plasma by means of generation of solitonlike structures moving into the plasma. In Fig. 1, the temporal evolution of a laser pulse in the form of $a_i = a_o \exp[-(t - 2\tau)^4/2\tau^4]$ with $a_o = 0.71$, $\tau = 100$, and interacting with a plasma with $n_o = 1.3$ is shown. At these pulse parameters, two solitary single-humped waves are generated at the plasma boundary and then slowly propagate as quasistationary plasma-field structures with a velocity ($v_s \approx 0.2c$) substantially lower than the speed of light. It should be emphasized that such single-humped structures (without nodes and having specific properties corresponding to a discrete spectrum of propagation velocities [13]) theoretically should exist only in slightly overdense plasmas where $n_o - 1 \ll 1$ and for weakly relativistic amplitudes; these are pure relativistic solitons where the contribution to the refraction index due to electron density perturbation is much weaker than the one due to the relativistic nonlinearity. However, it is important to note that our simulations show the significant role played by such single-humped structures also in the penetration dynamics for higher overdense plasma densities and higher incident amplitudes. Since the field distribution of these structures is quite close to the pure relativistic localized solutions (without density perturbations), the incident pulse duration must exceed some critical value in order to be able to excite such solitons. The critical pulse length can be estimated as $\tau^* \sim l_s/v_s$, where l_s is the width of the structure. For the incident amplitude presented in Fig. 1, each generated soliton requires a time $\sim\tau^* \approx 50$, which is in good agreement with the

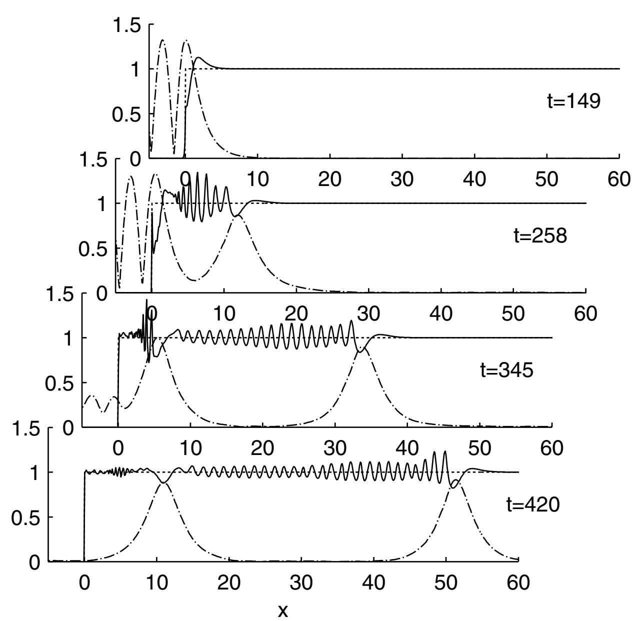

FIG. 1. Snapshots of the time evolution of the electron density (continuous line) and the solitary structures (dash-dotted line) generated by a laser pulse with $a_o = 0.71$, $\tau = 100$, and propagating into a plasma with $n_o = 1.3$ and a corresponding threshold value for penetration equal to $a_{\rm th} = 0.62$.

simulation. The generation of similar structures can also be inferred from the results of PIC simulations, such as those presented in [6]. It should also be emphasized that, at moderate relativistic intensities, this regime is rather insensitive to the ion mobility due to the comparatively fast motion of these structures.

At higher background densities, $n_o > 1.5$, the dynamic regime of interaction is completely different, as shown in Fig. 2, where a laser pulse with $a_o = 1.9$ interacts with a plasma with $N_o = 1.6N_c$. The earliest stage of the spatial evolution presents the characteristic distribution of a nonlinear skin layer, but the ponderomotive force acting at the vacuum-plasma boundary is pushing electrons into the plasma, thus shifting the real boundary to a new position. When the field amplitude on the real boundary exceeds the threshold, the interaction leads to the creation of a deep electron cavity. The whole plasma-field structure then starts to penetrate into the plasma and the same process is later repeated at the boundary. However, the electromagnetic field penetrates only a finite length into the plasma (as seen in Fig. 2 to a depth of $x \approx 10$). What is interesting is that, for comparatively long pulses, after a transient stage, the plasma settles down into a quasistationary plasma-field distribution, allowing for penetration of the laser energy over a finite length only, a length which increases with increasing incident intensities. The electron density distribution becomes structured as a sequence of electron layers over the ion background, separated by about half a wavelength with depleted regions. The peak electron density increases from layer to layer reaching an absolute maximum in the layer closest to the vacuum boundary. At the same time, the width of the layers becomes more and more narrow. Such nonlinear plasma structures act as a distributed Bragg reflector and they are quite close to those described analytically in [14]. However, as these structures during the interaction dynamics move very slowly or even are motionless, the concomitant strong charge separation together with the finite ion mobility may have a strong impact on the dynamics.

On one hand, it is evident that the cavitation structures cannot be produced in a electron-positron plasma ($\delta = 1$) since the ponderomotive force acts equally on negative and positive charges. In this case the interaction dynamics (as also follows from our simulations) is more simple and results in the formation of a plasma shock wave, without charge separation, that propagates into the plasma. On the other hand, it is easy to estimate the time during which the ions may still be considered as immobile in the cavities: This time may be taken as $t^* \approx L/v_i$, where L is the width of the cavity and v_i is the ion velocity. For motionless structures with $L \approx \lambda/2$ and completely uncompensated ion charge (see Fig. 2), we have $t \approx \omega_i^{-1}$. This means that, for pulse durations less than ω_i^{-1}, penetration of laser energy into overdense plasmas must occur through excitation of such deep cavities. After this time, ions will start to move resulting in the formation of narrow ion layers located around the electron layers; thus, in such interaction dynamics, the structural lamination of the plasma will be followed by the formation of a shock wave which penetrates into the plasma. We also expect that such penetration dynamics can be important even for longer pulses but for higher incident intensities. This is due to the fact that during the transient stage these structures are moving with velocities increasing with the incident intensities, and therefore they can avoid significant influence of ion motion by avoiding forming an ion inhomogeneity. This is clearly seen in Fig. 3, where a laser pulse with $a_o = 2.6$,

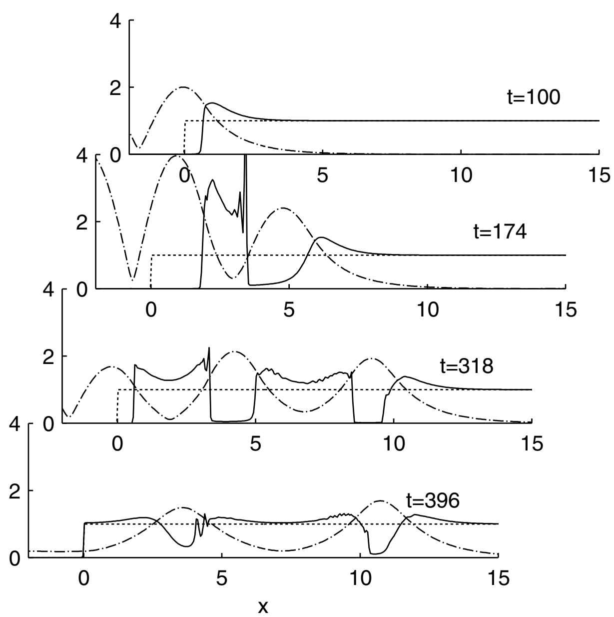

FIG. 2. Snapshots of the time evolution of the electron density (continuous line) and the field (dash-dotted line) for $a_o = 1.9$, $n_o = 1.6$, and $a_{\mathrm{th}} = 0.99$; other parameters as in Fig. 1.

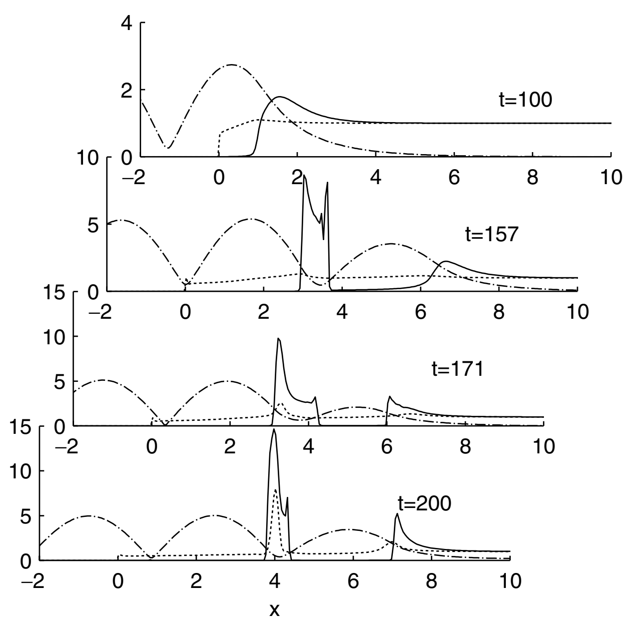

FIG. 3. Snapshots as in Fig. 2 but including ion dynamics (dashed line) for $a_o = 2.6$, $\delta = 1840$.

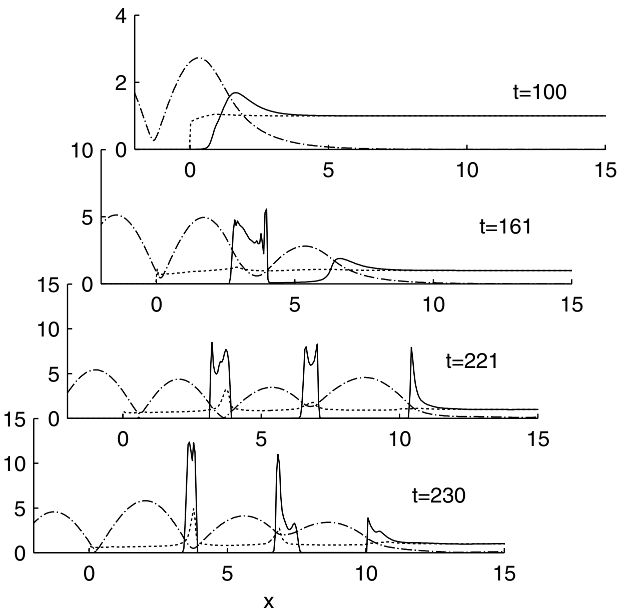

FIG. 4. The same as in Fig. 3 but for $\delta = 3860$.

$\tau = 100$, an intensity of 1.44×10^{19} W/cm^2, a wavelength of 1 μm, and 100 fs long interacts with a hydrogen ($\delta = 1840$) plasma with the same density as in Fig. 2, $N_o = 1.6N_c$. For these parameters $t^* \approx 40$ fs. Initially, penetration occurs on a smooth ion background, i.e., as in the immobile ion model in Fig. 2, up to a time $t \approx 170$. Only after this, ions start to produce an inhomogeneity in the form of an ion layer at $x \sim 4$ and an ion peak at $x \sim 7$ at the time $t = 200$. This feature will later be transformed into a plasma shock wave which slowly compresses the plasma. It is important to note that the effective ponderomotive force producing the shock wave corresponds to intensity values less than the maximum incident one and usually is of the order of the threshold of SIT. In the case with mobile ions, the self-induced transparency regime is realized at higher incident intensities than for fixed ions (in Fig. 2, $a_o = 1.9$). At lower amplitudes $a_i < 2.6$, the interaction dynamics is completely different and occurs through formation of plasma shock waves with low levels of charge separation. For the case of heavy ions, which is important for possible applications, deposition of laser energy in the plasma through the excitation of such cavitation structures in the SIT regime will be illustrated more clearly. In Fig. 4 we present results for the same parameters as in Fig. 3, but for $\delta = 3680$, roughly describing completely stripped multicharge heavy ions that can occur at relativistic intensities. As is clearly seen, the relativistic laser field with a pulse duration of 100 fs penetrates into the overdense plasma, producing multielectron layers separated by pure ion regions where strong electrostatic fields due to charge separation are produced. Obviously, these fields are responsible for fast ion beam generation which is not included in our fluid model. However, since the ion beams absorb only a small amount of the laser energy, we can expect that the macrodynamics of the interaction will exhibit the same features; moreover, by using these generated plasma-field structures we can estimate possible energetic parameters of the fast ions. On the longer time scale, the ion density distribution will shrink and will be localized around the electron layers (at $x \sim 4, 7$) and only the last layer (at $x \sim 10$ at the moment $t = 230$) will form a plasma shock wave.

In conclusion, we have shown that, depending on the background supercritical density, there are two qualitatively different scenarios of laser energy penetration into overdense plasmas in the regime of relativistic self-induced transparency. For plasma densities less than $N_o < 1.5N_c$, the penetration of the laser energy occurs in the form of long-lived solitonlike structures which are generated at the vacuum-plasma boundary plasma and then propagate into the plasma with low velocity. At higher densities, the interaction can result in the generation of plasma-field structures consisting of alternating electron and subsequent ion layers. The electromagnetic energy then penetrates into the overdense plasma over a finite length only, as determined by the incident intensity.

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