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Instability of the overdense plasma boundary induced by the action of a powerful photon beam

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An analytic solution describing an instability of a vacuum–overdense-plasma interface under the action of a powerful laser pulse is developed. The explicit dispersion relations and the space and time dependence of the perturbations have been obtained in the linear approximation. The influence of the instability on the interaction process is discussed and some comparisons with previous studies are presented.

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I. INTRODUCTION

The rapid advances in laser science and technology during the past decade have made possible the development of high-peak-intensity, subpicosecond, high-contrast-ratio ($\sim 10^{10}$) lasers. With such lasers on-target flux densities in excess of $10^{19}-10^{20}$ W/cm^2 for pulse duration as short as 10^{-13} sec will soon be available in several laboratories [1]. Simple estimates show that the ponderomotive forces of the electromagnetic field for such intensities may be of major importance in laser-matter interactions. The inward-directed force acting on the electrons appears to be larger than the thermal pressure gradients responsible for expansion. Hence, a very short time after the beginning of the interaction ($t \gg \omega_{pe}^{-1}$, where ω_{pe} is the electron plasma frequency), the ions start an inward motion dragged by the ambipolar field. Consequently, the plasma-vacuum interface acquires an inward acceleration under the action of the ponderomotive force of the electromagnetic field.

Acceleration of the plasma-vacuum boundary, which has a steplike density gradient, by the "light photon fluid" produces a clear example of Rayleigh-Taylor-like instability induced by the electromagnetic source. The Rayleigh-Taylor instability, which has been studied for a long time in laser-fusion plasmas, manifests itself near the ablation surface, where the gradients of the thermal pressure and density have opposite directions [2]. The ponderomotive force in the "corona" of a laser-produced plasma becomes significant in the underdense ($n_e \leq n_c$, where n_c is the electron critical density) even at relatively low intensities $I \geq 10^{14}$ W/cm^2 in accordance with Ref. [3]. The hydrodynamic instability excited in the underdense corona ($n_e \leq n_c$) and the resulting turbulence ("bubble" formation) in laser fusion conditions ($I \sim 10^{14}$ W/cm^2) have been observed in the numerical calculations of Ref. [4]. It is worth noting that in the conditions considered in that reference the interplay of many processes (radiation trapping, various forms of filamentation, etc.) was of importance along with the ponderomotive force of the incident wave. Recently [5], bubble formation was observed in particle-in-cell (PIC) simulations of laser-plasma interactions using very short ($\omega_0 t \sim 200$) and powerful ($I\lambda^2 = 10^{19}$ W μm^2/cm^2) laser pulses, where the ponderomotive force of the incident beam dominated the interaction. The plasma created in such an interaction is transient, nonequilibrium, and nonlocal [5–7].

The ponderomotive force at the boundary of such a plasma has been calculated recently [8]. The aim of this paper is to solve analytically the problem of stability of the steplike, overdense, nonlocal plasma-vacuum boundary accelerated inwards by the ponderomotive force. As is usual for this kind of problem, we will study the stability of the initial state of equilibrium, when the ponderomotive force supports the plasma boundary. The character of the equilibrium of this initial static state will be determined by introducing small initial perturbations to the various quantities and by following their evolution. The plan of this paper is as follows. In Sec. II the problem is formulated. The character, the form, and relations between the initial perturbations are described in Sec. III. The equations for the perturbations and the boundary conditions are derived in Sec. IV. The explicit solutions and dispersion relation are obtained in Sec. V. In Sec. VI we discuss the results and possible influence of the instability on the interaction process, compare the results with previous studies, and estimate the following turbulence characteristics. We draw the conclusions in Sec. VII.

II. FORMULATION OF THE PROBLEM

Let us suppose that a powerful plane-polarized laser beam is incident along the normal to a target (the z axis coincides with the direction of the normal) and creates on the target surface an energy flux density exceeding the relativistic value $I > I_r = 4 n_c m_e c^3 = 1.14 \times 10^{19} [\lambda(\mu\text{m})]^{-2}$ W/cm^2 (n_c, m_e, c, and λ are electron critical number density, electron mass, speed of light in vacuum, and wavelength of the incident light in micrometers, respectively). The incident-field components have the form

$$\mathbf{E}(E_{\text{in}},0,0); \mathbf{B}(0,B_{\text{in}},0) \sim \exp\{-i\omega_0 t + i k_0 z\} .$$

Here $k_0 = \omega_0 / c$. In accordance with Refs. [6,7], we assume that the beam ionizes the atoms in the target in a time shorter than a femtosecond due to both field (tunnel) and electron-impact ionization processes. Hence the laser beam interacts with a homogeneous, collisionless plasma which has a nonlocal relation between the current and the electric field inside the plasma. For the case con-

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sidered here the only force acting on the electrons in the z direction is the Lorentz force

$$\mathbf{F}=(1/c)[\mathbf{j}\times\mathbf{B}] \ . \tag{1}$$

The spatial dependence of the magnetic field, electric current, and ponderomotive force averaged over the laser period were calculated in Ref. [8]:

$$\begin{aligned} j_x(\eta)&=\frac{cB_0}{4\pi^2 l_s}\int_{-\infty}^{\infty}\frac{d\theta\exp(i\theta\eta)}{1+i|\theta|\theta^2}\ , \quad \eta=\frac{z}{l_s}\ , \\ B_y(\eta)&=B_0\left\{1-\frac{1}{i\pi}\int_0^{-\infty}\frac{d\theta\sin\theta\eta}{1+i\theta^3}\right\} . \end{aligned} \tag{2}$$

Here B_0 is the magnetic-field value on the vacuum-plasma boundary ($z=0$) and l_s is the field penetration depth into the plasma (the skin depth). For the case of low absorption, one can estimate B_0 as

$$B_0\approx 2B_{\rm in} \tag{3}$$

and relate it to the intensity of the incident light

$$I=cB_{\rm in}^2/4\pi \ . \tag{4}$$

Later in this paper we will use the value of the ponderomotive force on the boundary of the plasma ($z=0$) while solving the problem of the stability of the interface. This value can be obtained by evaluating the integrals in (2) along with the value of the electric current

$$\begin{aligned} j_x(z=0)&\equiv j_0=\frac{2}{3}\frac{cB_0}{4\pi l_s}\ , \\ F_z(z=0)&\equiv F_0=\frac{2}{3}\frac{B_0^2}{4\pi l_s}\ . \end{aligned} \tag{5}$$

For further simplification let us use for the skin depth the formula for the stationary skin effect, thus neglecting its slow dependence with time ($t\gg\omega_0^{-1}$) and incident intensity, which are characteristic of the nonstationary anomalous skin effect [6,7]. Hence we assume that

$$l_s\approx c/\omega_{pe}\sim n_e^{-1/2} \tag{6}$$

and, consequently, $F_0\sim j_0\sim n_e^{-1/2}$, thus preserving the density dependence of the force, which is most important for studying the problem of the hydrodynamic instability. We make also the usual assumptions for linear stability problems in magnetohydrodynamics, namely, quasineutrality, incompressibility, and assume the plasma to be a perfect conductor. Thus the full set of equations for the problem can be written as follows:

$$\begin{aligned} &\rho\left[\frac{\partial\mathbf{v}}{\partial t}+(\mathbf{v}\nabla)\mathbf{v}\right]=-\mathbf{F}p-\frac{\partial p}{\partial\mathbf{x}}\ , \\ &\mathrm{rot}\mathbf{E}=-\frac{1}{c}\frac{\partial\mathbf{B}}{\partial t}\ , \\ &\mathrm{rot}\mathbf{B}=\frac{4\pi}{c}\mathbf{j}\ , \\ &\mathrm{div}\mathbf{B}=0\ , \quad \mathrm{div}\mathbf{v}=0\ , \\ &\frac{\partial\rho}{\partial t}+(\mathbf{v}\nabla)\rho=0\ . \end{aligned} \tag{7}$$

We neglected the displacement current in (7) in accordance with the condition

$$\rho\gg B_0^2/4\pi c^2 \quad \text{or} \quad c\gg V_A=B_0/(4\pi\rho)^{1/2}\ ,$$

which is valid for a dense plasma ($\omega_{pe}\gg\omega_0$, where ω_{pe} is plasma frequency) with nonrelativistic ion motion and V_A is the Alfvén velocity. As follows from the first equation of (7), the initial state of equilibrium corresponds to the situation when the ponderomotive force is balanced by the thermal pressure gradient in the matter.

III. INITIAL PERTURBATIONS

Let us introduce small perturbations to all initial quantities, namely, velocity, pressure, mass density, current, and electric and magnetic fields: $\mathbf{u}$, δp, $\rho+\delta\rho$, $\mathbf{j}_0+\delta\mathbf{j}$, $\mathbf{E}+\mathbf{e}$, $\mathbf{B}_0+\mathbf{h}$. In the problem being considered, the electrons start to move instantaneously, dragging the ions after time of the order of ω_{pe}^{-1}. Hence the initial perturbations are related to the initial inhomogeneities as well as to the perturbations created by the initial electron motion (electron seeded perturbations [5]). It follows from the previous studies of the flute, or interchange, instabilities in a low-density plasma [9], that the hydrodynamic response is of major importance even for a rare plasma. We will consider the influence of the electron motion on this instability elsewhere. Here we consider only hydrodynamic instabilities, assuming the ions respond immediately to the ponderomotive force, and we thus neglect any fast electron oscillations. The development of the electromagnetic perturbations is related to the plasma motion (10). Hence

$$\begin{aligned} \mathbf{e}&=-\frac{1}{c}(\mathbf{u}\times\mathbf{B}_0)\ , \\ \frac{\partial\mathbf{h}}{\partial t}&=\mathrm{rot}(\mathbf{u}\times\mathbf{B}_0)\ . \end{aligned} \tag{8}$$

The perturbation of the current is related to both the perturbations of the density and velocity:

$$\delta\mathbf{j}=\left[\frac{\partial\mathbf{j}}{\partial\rho}\right]_u\delta\rho+\left[\frac{\partial\mathbf{j}}{\partial\mathbf{u}}\right]\delta\mathbf{u}\ .$$

Making use of (5), (6), and (8), one obtains

$$\delta\mathbf{j}=\frac{c}{4\pi}\mathrm{rot}\mathbf{h}+\mathbf{j}_0\frac{\delta\rho}{2\rho}\ . \tag{9}$$

Taking into account the identities of vector analysis and the solenoidal nature of the velocity and the magnetic field, one can reduce the equation for the magnetic-field perturbation to a simpler form,

$$\frac{\partial\mathbf{h}}{\partial t}=(\mathbf{B}_0\nabla)\mathbf{u}-(\mathbf{u}\nabla)\mathbf{B}_0\ . \tag{10}$$

For the sake of simplicity we will neglect in the further analysis the space derivative of the initial magnetic field in Eq. (10). This term does not change the final form of the dispersion relation but makes the equation for the space dependence of the perturbations more cumbersome.

IV. EQUATIONS FOR THE PERTURBATIONS AND BOUNDARY CONDITIONS

One can now linearize the set of equations (7) taking into account (8)–(10):

$$\begin{aligned}
&\rho\frac{d\mathbf{u}}{dt}=\frac{1}{c}(\mathbf{j}_0\times\mathbf{h})-\frac{1}{c}(\delta\mathbf{j}\times\mathbf{B}_0)-\frac{\partial\delta p}{\partial\mathbf{x}}\ ,\\
&\frac{\partial\mathbf{h}}{\partial t}=(\mathbf{B}_0\nabla)\mathbf{u}\ ,\\
&\delta\mathbf{j}=\frac{c}{4\pi}\mathrm{rot}\mathbf{h}+\mathbf{j}_0\frac{\delta\rho}{\rho}\ ,\\
&\mathrm{div}\mathbf{h}=0\ ,\quad \mathrm{div}\mathbf{u}=0\ ,\\
&\frac{\partial\delta\rho}{\partial t}+\mathbf{u}\frac{\partial\rho}{\partial\mathbf{x}}=0\ .
\end{aligned}\tag{11}$$

Making use of $\mathbf{j}_0(j_0;0;0)$, $\mathbf{B}_0(0;B_0;0)$, and the dependence of density on the z coordinate only, one can reduce (11) to the form

$$\begin{aligned}
&\rho\frac{du_x}{dt}=\frac{B_0}{c}(\delta j)_z-\frac{\partial\delta p}{\partial x}\\
&\rho\frac{du_y}{dt}=\frac{1}{c}j_0h_z-\frac{\partial\delta p}{\partial y}\\
&\rho\frac{du_z}{dt}=-\frac{1}{c}j_0h_y-\frac{1}{c}B_0(\delta j)_x-\frac{\partial\delta p}{dz}\\
&\frac{\partial\mathbf{h}}{\partial t}=B_0\frac{\partial\mathbf{u}}{\partial y}\ ,\quad \delta\mathbf{j}=\frac{c}{4\pi}\mathrm{rot}\mathbf{h}+\mathbf{j}_0\frac{\delta\rho}{2\rho}\ ,\\
&\mathrm{div}\mathbf{h}=0\ ,\quad \mathrm{div}\mathbf{u}=0,\quad \frac{\partial\delta\rho}{\partial t}+u_z\frac{\partial\rho}{\partial z}=0\ .
\end{aligned}\tag{12}$$

We will look for solutions in the form of an expansion into the normal modes u, δp, $\delta\rho$, etc. [10]

$$\sim\int\exp\{\gamma t+ik_xx+ik_yy\}f(z)dk_xdk_y\ .\tag{13}$$

Let us reduce the set (12) to the equations for the velocity and pressure components only:

$$\begin{aligned}
&\gamma\rho u_x=\frac{B_0^2}{4\pi\gamma}(k^2u_x-ik_xDu_z)-ik_x\delta p\ ,\\
&\gamma\rho u_y=\frac{j_0B_0}{c\gamma}ik_yu_z-ik_y\delta p\ ,\\
&\gamma\rho u_z=\frac{j_0B_0}{c\gamma}(ik_xu_x+Du_z)\\
&\qquad-\frac{B_0}{c}\left\{\frac{cB_0}{4\pi\gamma}(-k_y^2u_z+ik_xDu_x+D^2u_z)-\frac{j_0D\rho}{2\gamma\rho}u_z\right\}\ ,\\
&\mathrm{div}\mathbf{u}=0\ .
\end{aligned}\tag{14}$$

We denote here $k^2=k_x^2+k_y^2$ and $D=d/dz$. Now eliminating u_x, u_y, and δp from (14) one can reduce (14) to an equation for u_z alone:

$$\frac{1}{\rho}D(\rho Du_z)-k_y^2\frac{B_0^2}{4\pi\gamma^2}D^2u_z=\left\{1-k_y^2\frac{B_0^2}{4\pi\rho\gamma^2}-\frac{j_0B_0}{\rho c\gamma^2}\frac{D\rho}{2\rho}-\frac{k_x^2k_y^2}{k^2\rho\gamma}\frac{\left[\dfrac{j_0B_0}{c\gamma}\right]^2}{\rho\gamma-k_y^2\dfrac{B_0^2}{4\pi\gamma}}\right\}k^2u_z\ .\tag{15}$$

Substituting into (15) the formula (5) for j_0 and denoting

$$\gamma_0^2=k\frac{j_0B_0}{2c\rho}\equiv\frac{kV_A^2}{3l_s}\ ,\quad V_A^2=\frac{B_0^2}{4\pi\rho}\ ,\tag{16}$$

Eq. (15) can be reduced to the form

$$\frac{\gamma^2}{\rho}D(\rho Du_z)-k_y^2V_A^2D^2u_z=u_zk^2\left\{\gamma^2-k_y^2V_A^2-\frac{\gamma_0^2(D\rho)}{k\rho}-\frac{4k_x^2k_y^2}{k^4}\frac{\gamma_0^4}{\gamma^2-k_y^2V_A^2}\right\}\ .\tag{17}$$

Here V_A is the familiar Alfvén velocity. Note that for the problem considered the magnetic field in (16) is the time-averaged magnetic field of the incident electromagnetic wave, while in the conventional definition of the Alfvén velocity B is a constant magnetic field applied to the plasma. The boundary conditions which must be satisfied at the vacuum-plasma interface are u_z and $\mathbf{h}$ and are both continuous. Due to the second equation of (11) only one condition of continuity of u_z remains. The second boundary condition follows from (17) by integrating it over the infinitesimal distance across the interface. We obtain

$$\{\gamma^2(\rho Du_z)\}_{z=-0}^{z=+0}-\{k_y^2V_A^2(Du_z)\}_{z=-0}^{z=+0}=-\gamma_0^2k(u_z)_{z=0}\ .\tag{18}$$

V. SOLUTIONS

In a plasma where the density is constant, Eq. (7) reduces to the simple form

$$D^2u_z=k^2\left[1-\frac{4k_x^2k_y^2}{k^4}\frac{\gamma_0^4}{\left[\gamma^2-k_y^2V_A^2\right]^2}\right]u_z\ . \tag{19}$$

This equation has the general solution

$$u_z=A\ \exp\{\pm qz\}\ , \tag{20}$$

where A is a constant and q is defined by

$$\frac{q^2}{k^2}=1-\frac{4k_x^2k_y^2}{k^4}\frac{\gamma_0^4}{(\gamma^2-k_y^2V_A^2)^2}\ . \tag{21}$$

The solution of (20) with a positive sign has to be used for the region $z<0$, while the solution with a negative sign corresponds to the region $z>0$ in order to obtain the perturbations vanishing at infinity. Introducing (20) into (18), we obtain the dispersion relation for the problem:

$$\gamma_0^2k=q\left[\gamma^2-k_y^2V_A^2\right)\ . \tag{22}$$

Combining (21) and (22) and solving for the growth rate of the instability one can obtain the explicit relation

$$\gamma^2=k_y^2V_A^2\pm\gamma_0^2\left[1+\frac{4k_x^2k_y^2}{k^4}\right]^{1/2}\ . \tag{23}$$

Substituting (23) into (21), it is also easy to obtain the explicit expression for q,

$$q^2=k^2\left[1+\frac{4k_x^2k_y^2}{k^4}\right]^{-1}\ . \tag{24}$$

Thus Eqs. (20)–(24) represent the full formal solution of the problem.

VI. DISCUSSION

There is a clear analogy between the case considered here and classical Rayleigh-Taylor instability of a heavy fluid accelerated by air pressure, or, of instability of the interface between a heavy fluid imposed over a light one in a gravitational field [10,11]. Let us consider the case when $k_y=0$. In this case the positive sign in (23) corresponds to an instability with a growth rate

$$\gamma=(k_xV_A^2/3l_s)^{1/2}\ . \tag{25}$$

Hence the quantity $V_A^2/3l_s$ plays the role of the acceleration of gravity in our problem. Note that the other branch of the dispersion curve (with the positive sign) corresponds to surface waves propagating in the direction perpendicular to the lines of magnetic force. The other important property of the growth rate that has been obtained is its strong density dependence for long-wavelength perturbations along the magnetic field ($k_yl_s\ll1$). In this case, it follows from (25) that

$$\gamma^2\sim\rho^{-1/2}\ .$$

When $k_x=0$, we obtain from (23)

$$\gamma_{k_x=0}^2=\frac{k_yV_A^2}{3l_s}(3l_sk_y\pm1)\ . \tag{26}$$

In the long-wavelength limit ($3k_yl_s\ll1$) the dispersion relation (26) has two branches. One branch corresponds to unstable modes (positive sign), and the second one relates to waves propagating along magnetic force lines (similar to Alfvén waves). In the short-wavelength limit ($3k_yl_s\gg1$), there is only unstable branch. The growth rate in this case is density independent and has the form similar to the relation for the frequency of Alfvén waves:

$$\gamma^2\sim k_y^2V_A^2\ .$$

It is instructive to compare the solutions that have been obtained to the case when a constant magnetic field was applied parallel to an unstable interface of two fluids in a gravitational field [10,12]. In this case the magnetic force acts as surface tension, associated with the tension along the magnetic force lines. On the contrary, in the case considered in this paper all components of the ponderomotive force lead to an increase of the growth rate of the instability.

To understand the difference, let us compare the driving forces for both cases. In the case of the fluids in the gravity field, the gravitational force is responsible for the instability and the resulting fluid motion. The magnetic force, and the associated tension, appears only after the beginning of the fluid motion and acts in a direction opposite to the gravitational force, thus leading to stabilization. In the problem considered here the ponderomotive force (which has initially only the term related to the gradient of the energy density) supports the plasma in the initial equilibrium state, and is also the source of motion and the instability. The term associated with the tension along the magnetic force lines, appears due to the development of the perturbations and has the same direction as the unperturbed ponderomotive force. This is easy to understand, from a formal point of view if one expresses the ponderomotive force in terms of the magnetic field alone [10]

$$F_z=\frac{1}{c}[\mathbf{j}\times\mathbf{B}]_z=-\frac{\partial}{\partial z}\frac{|B|^2}{8\pi}+\frac{\partial}{\partial x_i}(B_zB_i)\ . \tag{27}$$

The last term in (27) is associated with the tension along the magnetic force lines. Note that (27) is valid only in the case when it is possible to neglect the displacement current.

Let us estimate the characteristic value of the growth rate of the instability γ_0 for the case when the wavelength of the perturbation equals the skin depth. For this case the growth rate is density independent. Thus, inserting $k=2\pi/l_s$ and $l_s=c/\omega_{pe}$ ($\omega_{pe}=4\pi e^2n_e/m_e$) into (25), one can obtain

$$\gamma_0^2=\frac{32\pi^2}{3}\frac{e^2Z_i}{m_eM_i}\frac{I}{c^3}\ . \tag{28}$$

Here Z_i, M_i, e, and m_e are ion and electron charges and masses, respectively, $I=cB_{\rm in}^2/4\pi$ is the intensity of the incident beam, and c is the speed of light in vacuum. Now we evaluate (28) for the conditions of the PIC simulations of Ref. [5], where an $I[\lambda(\mu\mathrm{m})]^2=1.2\times10^{19}$ W $\mu\mathrm{m}^2/\mathrm{cm}^2$ laser beam interacted with a plasma ($M_i/m_e=1836$) having an initial density profile in the form of a linear rise from 0 to $4n_c$ over a distance $4\lambda_0$ and a constant density for the next λ_0. For those conditions the instability growth time from (28) is $\gamma_0^{-1} \leq 4$ fs. It is easy to estimate that after a time of $(2-3)\gamma_0^{-1}$ the instability enters the nonlinear regime ($ak \geq 1$, where a is an amplitude of the perturbation), and turbulence will result. It follows from the two-dimensional calculations of Ref. 5 that at a time $\omega_0 t = 225$ ($t > 100$ fs) well-developed turbulence exists with a characteristic bubble (turbulent eddy) size of $10c/\omega_0$ ($\sim 1.5\lambda_0$). Hence the simple theory of this paper suggests, in agreement with the numerical simulations of Ref. [5], that the interaction of a very intense laser with a low-density plasma (for example, with foam targets, or a plasma with preformed-density gradient) enters the interaction regime with the turbulent plasma after several tens of femtoseconds.

Let us examine the ponderomotive force induced instability and the resulting turbulence from a more general point of view. For this purpose we present the Euler equation in the form for a velocity vortex rot**v**,

$$\frac{\partial}{\partial t}\mathrm{rot}\mathbf{v} - \mathrm{rot}(\mathbf{v} \times \mathrm{rot}\mathbf{v}) = -\frac{1}{4\pi\rho}\mathrm{rot}(\mathrm{rot}\mathbf{B} \times \mathbf{B}) \ .$$

One can draw at least two conclusions from this equation. First, in the quasistationary case, when the vortices (bubbles) are simply flowing with the inward plasma motion, the space scale of the vorticity depends only on the space scale of the driving magnetic field. Second, the fluid velocity is of the order of the Alfvén velocity [10] (note that for the case considered in this paper the magnetic field in the definition of the Alfvén velocity is the magnetic field of the incident wave averaged by the laser period). Hence the characteristic scale of the vortex has to be of the order of magnitude of the initial field scale (e.g., the skin depth).

VIII. CONCLUSION

To conclude, we have presented in this paper analytical solutions for the problem of the instability of a vacuum–overdense-plasma interface driven by the ponderomotive force of a powerful laser beam. The explicit formulas for the instability growth rate and spatial dependence of the perturbations are obtained in the linear approximation. The main features of the growth rate of the instability which are derived are its dependence on plasma density along with its strong dependence on the intensity of the incident beam. This instability is similar to the familiar Rayleigh-Taylor instability, where the term $V_A^2/3l_s$ [$V_A = B_{\mathrm{in}}/(4\pi\rho)^{1/2}$ is the Alfvén velocity and l_s is the field penetration depth] plays the same role as the acceleration due to gravity. One can also consider this instability as an example of unstable (with the growing amplitude) Alfvén waves, driven by the ponderomotive force of the incident laser beam. The characteristic time for the instability growth, γ^{-1}, appears to be of the order of several femtoseconds for an incident flux density of the laser beam in excess of 10^{18} W/cm^2. Thus one can expect the transition time to the nonlinear and turbulent regime of the plasma motion to be of the order of tens of femtoseconds. Hence the theory presented suggests that the intense ($I > 10^{18}$ W/cm^2), short (~ 100 fs) laser-beam–low-density-plasma interaction proceeds for most of the interaction time in the regime of turbulent plasma motion. The development of further theory will take into account the density and field gradients, as well as the initial electron motion.

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