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Interaction of a non-symmetric powerful laser pulse with a plasma

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The interaction of a powerful non-symmetric laser pulse with a plasma is studied. The non-symmetry is manifested in an abrupt cut-off of the rear edge of the laser pulse compared with its leading edge. At the same time, three qualitatively different regions are distinguished: the leading edge, the rear edge and the region behind the pulse, where it leaves a wake in the form of generated fields. An analytical solution has been found that defines the longitudinal accelerating field at the end of the rear edge. The results of numerical calculations confirm our physical point of view that the non-symmetry of the laser pulse increases the duration of the ion channel behind the front, thereby enhancing the focusing and effective acceleration of electron bunches.

The investigation of the interaction of a powerful laser pulse with a plasma and the phenomena accompanying this interaction is very important from the point of view of applications, in addition to its great scientific significance. This is due to the fact that the results of such investigations can be used directly for the solution of problems such as the radiation-frequency conversion (Wilks *et al.* 1988, 1989; Gildenburg *et al.* 1990), the generation of harmonics (Sprangle *et al.* 1990; Esarey *et al.* 1991; Esarey and Sprangle 1992), inertial-confinement fusion (Liu and Umstadter 1992; Tabak *et al.* 1994), excitation of intense relativistic longitudinal fields for the acceleration of charged particles (Tajima and Dawson 1979; Kaw *et al.* 1992; Andreev *et al.* 1992; Sprangle *et al.* 1992; Wurtele 1993; Borisov *et al.* 1992; Krall *et al* 1993; Esarey *et al.* 1993) and the initiation of nuclear effects (Boyer *et al.* 1988). Along with the generation of an intense relativistic longitudinal field by a powerful laser pulse, the search for the possibilities of effective extension of this field beyond the pulse front is of no less importance. In Berezhiani and Murusidze (1990, 1992) and Murusidze and Tsintsadze (1992) a strong dependence of the spatial–temporal structure of the generated longitudinal field on the shape of the laser pulse introduced into plasma was established. Garuchava and Tagviashvili (1994) proposed the concept of a non-symmetric laser pulse, the rear edge of which is cut off much more sharply than the leading edge (Fig. 1). With such a pulse structure, three qualitatively different regions are naturally distinguished: a leading edge, a rear edge and a region behind the pulse, where it leaves a wake in the form of

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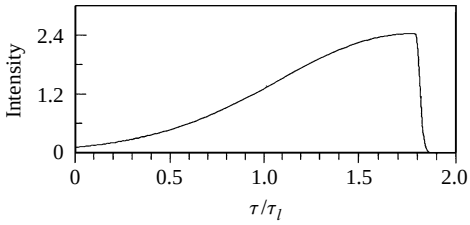


Figure 1. Profile of the normalized intensity $I = (e|E|^2/mc\omega_0)^2$ of an asymmetric pulse with Gaussian rise and fall; $(\tau_l/\tau_r)^2 = 10^3$.

generated fields and a positively charged channel. The formation of the channel as a result of forcing out of electrons by a high-frequency (HF) field localized in the pulse was predicted first by (Garuchava *et al.* (1986, 1990). In the region of the abrupt rear edge the longitudinal field and its trace behind the pulse should be more effective owing to the short time interval and the strong spatial gradients. The duration of the positively charged channel should be increased as well. These properties are confirmed by numerical calculations. If the pulse is symmetric and the duration of each half of the pulse is much longer than the period of electron plasma oscillations then the ion channel disappears at the end of the pulse, and the wake field following the pulse will not be effective for acceleration (Barness *et al.* 1987; Borisov *et al.* 1992; Mori *et al.* 1988; Sprangle *et al.* 1987; Sun *et al.* 1987; Tsintsadze *et al.* 1990). Here we concentrate our attention on the attractive properties of a powerful asymmetric laser pulse for the generation of an intense relativistical longitudinal field, describe the spatial and temporal structure of such a pulse, and determine analytically the maximum momentum of electrons at the end of the rear edge.

The duration of the leading front τ_l is much larger than the period of Langmuir oscillations, $\omega_p \tau_l \gg 1$, while for the duration of rear front τ_r the reverse equality holds, $\omega_p \tau_r \ll 1$. For the transverse scale $l_\perp$ of characteristic variations of physical values on the leading and rear edges, we have

$$\lambda_0 \ll l_{\perp l} \ll c\tau_l, \qquad \lambda_0 \ll c\tau_r \ll l_{\perp r}, \tag{1}$$

where $\lambda_0 = 2\pi k_0^{-1}$ is the wavelength of the HF wave. The period of this wave is much less than the durations of either the leading or rear edges, $\omega_0 \tau_l \gg 1$ and $\omega_0 \tau_r \gg 1$. Consequently, on both edges two time scales are distinguished. Corresponding to this, any of the values $S = (\mathbf{P}, n, \Phi, \mathbf{E}, \gamma)$ used in our problem can be represented as a sum of slowly and rapidly varying terms

$$S = \langle S \rangle + \tilde{S}, \tag{2}$$

where the angular brackets $\langle \ldots \rangle$ denote averaging over the rapid time scale $2\pi/\omega_0$. Here $\mathbf{P}$ is the relativistic momentum of electrons, n is their concentration, Φ is the potential of the generated longitudinal field, $\mathbf{E}$ is the electric field vector of the HF wave and γ is the relativistic factor. All characteristic velocities of our problem are much larger than the electron thermal velocity; thus hydrodynamic approximation is quite justified for the description of the interaction of a powerful laser pulse with a plasma. We also assume that all characteristic frequencies are much larger than the ion plasma frequency, and that the ion can be considered to be immobile. The corresponding set of

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equations includes the equation of motion, the equation of continuity and Maxwell's equations. It is convenient to introduce dimensionless quantities:

$$\frac{\mathbf{P}}{mc}\to\mathbf{P},\qquad \frac{e\Phi}{mc^2}\to\Phi,\qquad \frac{n}{n_0}\to n,\qquad \frac{e\mathbf{E}}{mc^2k_0}\to\mathbf{E}, \tag{3}$$

where n_0 is the equilibrium concentration of electrons. All dynamic values in the initial set of equations are represented in the form (2). Then, for slowly varying values, we obtain the following equations (Garuchava *et al.* 1986, 1990; Sprangle *et al.* 1992):

$$\frac{1}{c^2}\frac{\partial^2\langle\mathbf{P}\rangle}{\partial t^2}-\nabla^2\langle\mathbf{P}^t\rangle+\frac{1}{c}\frac{\partial}{\partial t}\nabla\langle\gamma\rangle+\frac{\omega_p^2}{c^2}\langle\mathbf{j}\rangle=0, \tag{4}$$

$$\frac{1}{c}\frac{\partial\langle n\rangle}{\partial t}+\nabla\cdot\langle\mathbf{j}\rangle=0, \tag{5}$$

$$\frac{1}{c}\frac{\partial\langle\mathbf{P}^l\rangle}{\partial t}=\nabla\langle\Phi\rangle-\nabla\langle\gamma\rangle, \tag{6}$$

$$\nabla^2\langle\Phi\rangle=\frac{\omega_p^2}{c^2}(\langle n\rangle-1), \tag{7}$$

$$\omega_p^2=\frac{4\pi e^2n_0}{m},$$

where the superscripts l and t indicate the irrotational and solenoidal components of a vector, and the nonlinear current is given by

$$\langle\mathbf{j}\rangle=\left\langle\frac{n\mathbf{P}}{\gamma}\right\rangle=\frac{\langle n\rangle\langle\mathbf{P}\rangle}{\langle\gamma\rangle}-\frac{\langle\tilde{n}\tilde{\mathbf{P}}\rangle}{\langle\gamma\rangle}-\frac{\langle n\rangle\langle\tilde{\mathbf{P}}\tilde{\gamma}\rangle}{\langle\gamma\rangle^2}-\frac{\langle\mathbf{P}\rangle\langle\tilde{n}\tilde{\gamma}\rangle}{\langle\gamma\rangle^2}. \tag{8}$$

The rapidly varying momentum and concentration are determined from

$$\frac{1}{c^2}\frac{\partial^2\tilde{\mathbf{P}}}{\partial t^2}-\nabla^2\tilde{\mathbf{P}}^t+\frac{1}{c}\frac{\partial}{\partial t}\nabla\tilde{\gamma}+\frac{\omega_p^2}{c^2}\tilde{\mathbf{j}}=0, \tag{9}$$

$$\frac{1}{c}\frac{\partial\tilde{n}}{\partial t}+\nabla\cdot\tilde{\mathbf{j}}=0, \tag{10}$$

$$\tilde{n}=\frac{c}{\omega_p^2}\nabla^2\tilde{\gamma}+\frac{c}{\omega_p^2}\frac{\partial}{\partial t}\nabla\cdot\tilde{\mathbf{P}}, \tag{11}$$

where

$$\begin{aligned}\tilde{\mathbf{j}}=&\frac{1}{\langle\gamma\rangle}\left(\langle n\rangle\tilde{\mathbf{P}}+\langle\mathbf{P}\rangle\tilde{n}-\frac{1}{\langle\gamma\rangle}\langle n\rangle\langle\mathbf{P}\rangle\tilde{\gamma}\right)\\&+\frac{1}{\langle\gamma\rangle^2}[\langle\gamma\rangle(\tilde{n}\tilde{\mathbf{P}}-\langle\tilde{n}\tilde{\mathbf{P}}\rangle)-\langle n\rangle(\tilde{\mathbf{P}}\tilde{\gamma}-\langle\tilde{\mathbf{P}}\tilde{\gamma}\rangle)\\&-\langle\mathbf{P}\rangle(\tilde{\gamma}\tilde{n}-\langle\tilde{\gamma}\tilde{n}\rangle)-\tilde{n}\tilde{\mathbf{P}}\tilde{\gamma}].\end{aligned} \tag{12}$$

Writing down the relativistic factor in the form (2), we have

$$\langle\gamma\rangle=(1+\tilde{\mathbf{P}}^2+\langle\mathbf{P}\rangle^2)^{1/2}, \tag{13}$$

$$\tilde{\gamma}=\frac{\tilde{\mathbf{P}}\langle\mathbf{P}\rangle}{\langle\gamma\rangle}\ll\langle\gamma\rangle. \tag{14}$$

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Furthermore, the solenoidal component of the rapidly varying momentum will be considered to be circularly polarized:

$$\tilde{\mathbf{P}}^t_\perp = \tfrac{1}{2}(\mathbf{e}_r + i\mathbf{e}_\varphi) P(r,z,t)\exp(-i\omega_0 t + ik_0 z + i\varphi) + \text{c.c.}, \tag{15}$$

$$\omega_0 = k_0 c,$$

where $\mathbf{e}_r$ and $\mathbf{e}_\varphi$ are the unit vectors along the radial and azimuthal axes; φ is the azimuthal angle. The dependences of rapidly varying values on coordinates and time is determined, mainly, by the exponential factor in (15). Using (15) and the equations $\nabla\cdot\tilde{\mathbf{P}}^t = 0$ and $\nabla\times\tilde{\mathbf{P}}^l = 0$ gives the useful relations

$$\tilde{P}^t_z = \frac{1}{k_0}\frac{\partial \tilde{P}^t}{\partial r}, \qquad \tilde{P}^l_r = \frac{1}{k_0}\frac{\partial \tilde{P}^l_z}{\partial r}, \tag{16}$$

$$|\tilde{\mathbf{P}}^t_\perp|^2 = |P|^2, \qquad \langle \tilde{P}^t \tilde{P}^t_r\rangle = 0, \tag{17}$$

$$\langle \tilde{P}^t_r \tilde{P}^t_z\rangle = \frac{i}{4k_0}\left(P^*\frac{\partial P}{\partial r} - \text{c.c.}\right), \tag{18}$$

$$\langle \tilde{P}^t \tilde{P}^t_z\rangle = \frac{1}{4k_0}\frac{\partial |P|^2}{\partial r}, \tag{19}$$

$$\langle (\tilde{P}^t_z)^2\rangle = \frac{1}{2k_0^2}\frac{\partial P}{\partial r}\frac{\partial P^*}{\partial r}. \tag{20}$$

Using these relations, it is easy to ascertain that the term in square brackets in (12) responsible for the generation of the harmonics of the fundamental frequency ω_0 makes a contribution that does not exceed the value

$$\frac{\langle n\rangle}{\langle\gamma\rangle}\frac{\tilde{\mathbf{P}}\cdot\langle\mathbf{P}\rangle}{\langle\gamma\rangle}\tilde{\mathbf{P}},$$

namely

$$\tilde{\mathbf{P}}\frac{1}{k_0}(\tilde{\mathbf{P}}\cdot\nabla)\frac{\langle n\rangle}{\langle\gamma\rangle}.$$

Thus, within the scope of the problem under investigation, we neglect this term in comparison with the term in large parentheses. From (10), (11) and (14), taking the laser pulse into consideration, we have the following inequalities: $\tilde{P}^l_\perp \ll \tilde{P}^l_z \approx \tilde{P}^t_z \ll \tilde{P}^t_\perp$. Variation of the transverse component of the fast momentum in time is connected with $\tilde{\mathbf{E}}_\perp$ and $\tilde{\gamma}$ by the equation

$$\frac{1}{c}\frac{\partial \tilde{\mathbf{P}}_\perp}{\partial t} = -k_0\tilde{\mathbf{E}}_\perp - \nabla_\perp\tilde{\gamma}.$$

Using this equation, and taking into account that $\tilde{\mathbf{P}}^t_\perp \approx \tilde{\mathbf{P}}_\perp$ and $\tilde{\mathbf{E}}^t_\perp \approx \tilde{\mathbf{E}}_\perp$, from (9) we have

$$\frac{1}{c^2}\frac{\partial^2\tilde{\mathbf{E}}_\perp}{\partial t^2} - \nabla^2\tilde{\mathbf{E}}_\perp - \frac{\omega_p^2}{c^2\omega_0}\frac{\partial}{\partial t}\tilde{\mathbf{j}}_\perp = 0. \tag{21}$$

We describe the dynamics of slowly varying values with variables $\tau = t - z/c$ and $\xi = z$. Let us assume that $(1/c)\,\partial/\partial\tau \gg \partial/\partial\xi$. This assumption is obvious since the dependence on the coordinate ξ describes the slow variation of the

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laser pulse during its evolution as a whole (its self-focusing and distribution of the values in the channel). Taking (1) into account, the equation for the leading edge is reduced to the following system:

$$2ik_0\frac{\partial E}{\partial \xi}+\nabla_\perp^2 E-\frac{\omega_p^2}{c^2}\frac{\langle n\rangle}{\langle\gamma\rangle}E=0, \tag{22}$$

$$\frac{1}{r}\frac{\partial}{\partial r}\left(r\frac{\partial\langle\Phi\rangle}{\partial r}\right)=\frac{\omega_p^2}{c^2}(\langle n\rangle-1), \tag{23}$$

$$\langle n\rangle=1+\frac{c^2}{\omega_p^2}\frac{1}{r}\frac{\partial}{\partial r}r\frac{\partial}{\partial r}\langle\gamma\rangle\geqslant 0, \tag{24}$$

$$\langle\gamma\rangle=(1+|E|^2)^{1/2}. \tag{25}$$

Here E is the amplitude of the HF field, $E=-iP$ (see (15)), and $\langle n\rangle$ cannot be negative. The boundary and initial conditions are written as

$$|E|^2(r,\xi=0,\tau)=|E_0|^2\exp\left(-\frac{r^2}{a_0^2}-\frac{\tau^2}{\tau_l^2}\right),\qquad \Phi(r,\xi,\tau=-\infty)=0,$$

$$n(r,\xi,\tau=-\infty)=1,\qquad \frac{\partial|E|^2(r=0,\xi,\tau)}{\partial r}=0.$$

Now we simplify the equations for the rear-edge region. In the period of time τ_r, the averaged values $\langle S\rangle$ have no time to change in the radial direction; therefore we suppose that

$$\frac{\partial\langle S\rangle}{\partial r}\ll\frac{1}{c}\frac{\partial}{\partial\tau}\langle S\rangle. \tag{26}$$

The second inequality in (1) characterizing the structure of the rear edge allows us to have

$$\frac{\partial\langle S\rangle}{\partial z}\approx-\frac{1}{c}\frac{\partial}{\partial\tau}\langle S\rangle. \tag{27}$$

Using (26) and (27) with the help of the equations

$$\boldsymbol{\nabla}\cdot\langle\mathbf{P}^t\rangle=0\quad\text{and}\quad\boldsymbol{\nabla}\times\langle\mathbf{P}^l\rangle=0,$$

we can determine the hierarchy of relations between the components $\langle\mathbf{P}^l\rangle$ and $\langle\mathbf{P}^t\rangle$:

$$\langle P_r^t\rangle\approx\langle P_r^l\rangle\ll\langle\gamma\rangle,\qquad\langle P_z^t\rangle\ll\langle P_r^t\rangle. \tag{28}$$

At the same time, $\langle P_r^l\rangle$ varies insignificantly on the rear edge:

$$\langle P_r^l\rangle\approx\langle P_r^l\rangle_0, \tag{29}$$

where $\langle P_r^l\rangle_0$ is the value of $\langle P_r^l\rangle$ at the end of the leading edge. On the rear edge, in contrast to the leading one, we must retain the mixed derivative $\partial^2/\partial\xi\,\partial\tau$ in the equation for determination of the HF field amplitude. Under the condition (26), this equation has the form

$$2ik_0\frac{\partial E}{\partial\xi}-2\frac{\partial^2 E}{\partial\xi\,\partial\tau}-\frac{\omega_p^2}{c^2}\frac{\langle n\rangle}{\langle\gamma\rangle}E=0. \tag{30}$$

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Equations (5) and (6) contain terms of order $(\lambda_0/l_{\perp r})^2$. Neglecting these terms under the conditions (26) and (27), (5) and (6) can be written as

$$\frac{\partial\langle n\rangle}{\partial\tau}-\frac{\partial}{\partial\tau}\frac{\langle n\rangle\langle P_z^l\rangle}{\langle\gamma\rangle}=0, \tag{31}$$

$$\frac{\partial^2\langle\Phi\rangle}{\partial\tau^2}=\omega_p^2(\langle n\rangle-1), \tag{32}$$

and from (6) we obtain

$$\frac{\partial}{\partial\tau}\langle P_z^l\rangle=-\frac{\partial}{\partial\tau}\langle\Phi\rangle+\frac{\partial}{\partial\tau}\langle\gamma\rangle. \tag{33}$$

Consequently,

$$\langle n\rangle=\frac{C_1(r,\xi)}{1-\langle P_z^l\rangle/\langle\gamma\rangle}, \tag{34}$$

$$\langle P_z^l\rangle=-[\langle\Phi\rangle+C_2(\xi,r)]+\langle\gamma\rangle. \tag{35}$$

$C_1(r,\xi)$ and $C_2(r,\xi)$ are determined by the magnitudes of the quantities at the end of the leading edge:

$$C_1(r,\xi)=\langle n_0\rangle\left(1-\frac{\langle P_z^l\rangle_0}{\langle\gamma\rangle_0}\right). \tag{36}$$

$$C_2(r,\xi)=\langle\gamma\rangle_0-\langle P_z^l\rangle_0-\langle\Phi\rangle_0. \tag{37}$$

We should expect that $|\langle P_z^l\rangle|\geqslant 1$ on the rear edge. We confirm below that this is indeed true. Thus, for the relativistic factor on the rear edge, we have to write the following:

$$\langle\gamma\rangle=(1+|E|^2+\langle P_z^l\rangle^2)^{1/2} \tag{38}$$

From (35), we find

$$\langle P_z^l\rangle=\frac{1}{2}\left\{\frac{1+|E|^2}{C_2(\xi,r)+\Phi}-[C_2(r,\xi)+\Phi]\right\}. \tag{39}$$

According to (32), the variation of the potential on the rear edge

$$\delta\langle\Phi\rangle\approx(\omega_p\tau_r)^2\ll 1,$$

and its influence is negligible. As $\langle P_z^l\rangle_0\ll\langle\gamma\rangle_0$, from (36) we obtain

$$C_1(r,\xi)\approx\langle n\rangle_0, \tag{40}$$

and for $\langle P_z^l\rangle$ we have

$$\langle P_z^l\rangle=\frac{1}{2}\frac{|E|^2-|E_0|^2}{(1+|E_0|^2)^{1/2}}+\langle P_z^l\rangle_0. \tag{41}$$

Consequently, at the end of the rear edge where $|\mathbf{E}|=0$, $\langle P_z^l\rangle$ achieves the maximum value

$$|\langle P_z^l\rangle|_{\max}=\frac{1}{2}\frac{|E_0|^2}{(1+|E_0|^2)^{1/2}}\langle P_z^l\rangle_0. \tag{42}$$

For $|E_0|^2\gg 1$,

$$|\langle P_z^l\rangle|_{\max}\approx\tfrac{1}{2}|E_0|\gg\langle P_z^l\rangle_0. \tag{43}$$

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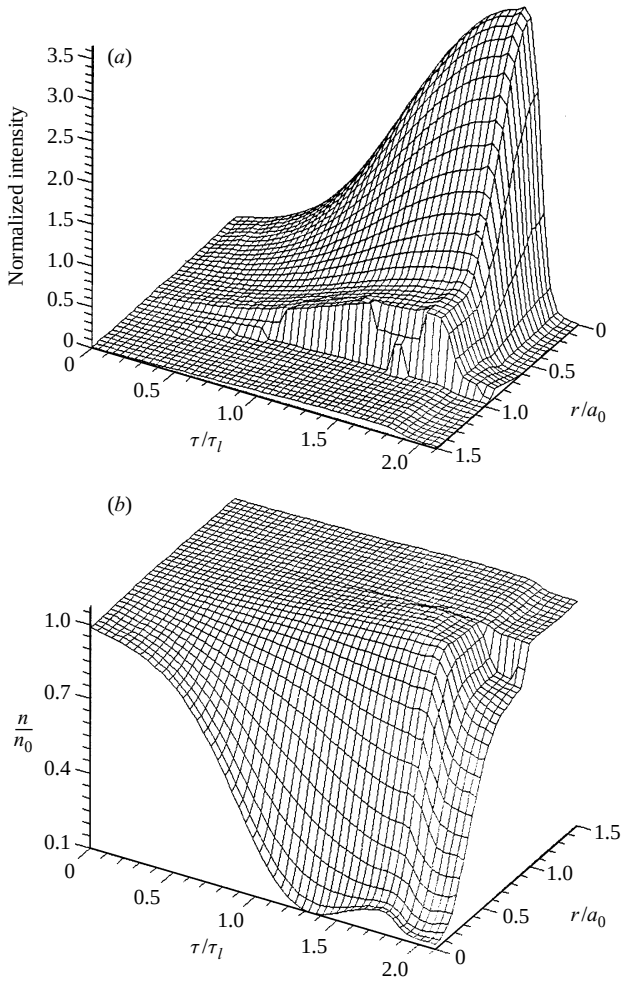


Figure 2. The normalized intensity distribution of the laser pulse (a) and the corresponding profile of the plasma electron density (b) after travelling a distance of $0.4L_R$ (where $L_R = \pi a\ 20/\lambda_0$ is the Rayleigh length). It shows one of the first stages of ion-channel formation.

We can now justify the relation (29). Under the conditions (26) and (27), the equation $\nabla \times \langle \mathbf{P} \rangle| \ = 0$ has the solution

$$\langle P_r^l \rangle \approx -\frac{c\tau_r}{l_{\perp r}} \langle \overline{P_z^l} \rangle + \langle P_r^l \rangle_0,$$

where $\langle \overline{P_z^l} \rangle$ is the average magnitude of $\langle P_z^l \rangle$ in the interval τ_r. Therefore the relation (29) suggests that

$$\frac{c\tau_r}{2l_{\perp r}} |E_0| \approx \langle P_r^l \rangle_0.$$

This equality is easily satisfied. The relation (29) has a clear physical meaning: $\langle P_r^l \rangle$ is determined by radial redistribution of electrons, as a result of which a positively charged ion channel is formed. The fact that $\langle P_r^l \rangle$ almost does not

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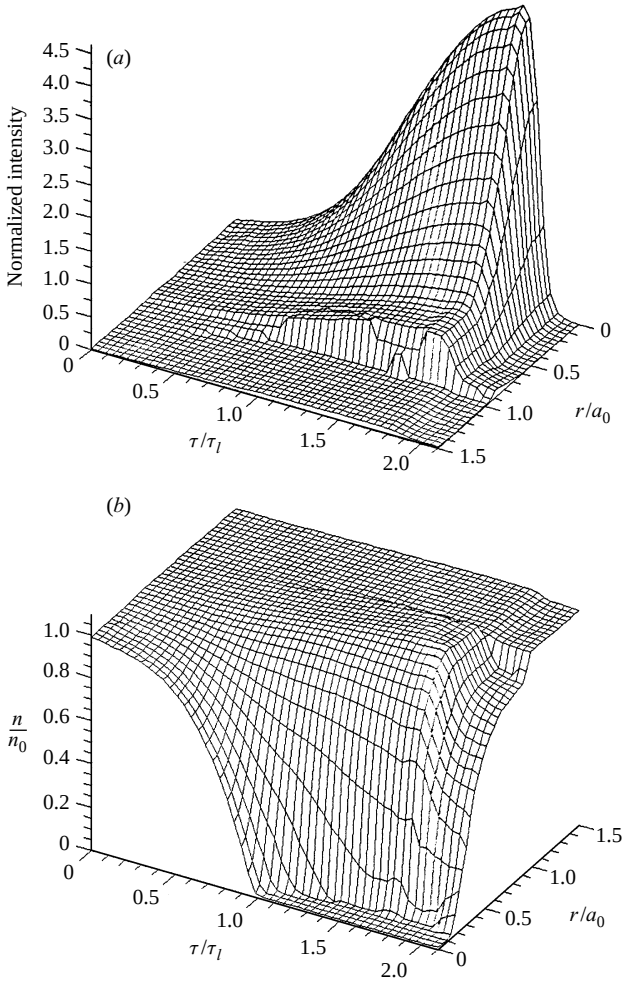


Figure 3. The intensity (*a*) and the corresponding electron-density profile (*b*) after the pulse has travelled a distance of $0.5L_R$. The ion channel is formed, and its further evolution continues, as seen in the subsequent figures.

vary on the rear edge means that the electrons have no time to be redistributed to any great extent in the time interval τ_r, and the ion channel formed by the leading edge also remains almost undistorted.

The numerical calculations were carried out for the following values of the laser-pulse and plasma parameters: $(a_0/\lambda_0)^2 = 4 \times 10^2$, $(\omega_p/\omega_0)^2 = 10^{-3}$, $(\tau_l/\tau_r)^2 = 10^3$ and $|E_0|^2 = 2.44$. Space–time coordinates were normalized as follows:

$$\xi \to \frac{\xi}{4L_R}, \qquad r \to \frac{r}{l_\perp}, \qquad \tau \to \frac{\tau}{\tau_l},$$

where

$$L_R = \pi \left(\frac{a_0}{\lambda_0}\right)^2 \lambda_0$$

is the Rayleigh length. Figures 2–6 make it possible to follow the processes of

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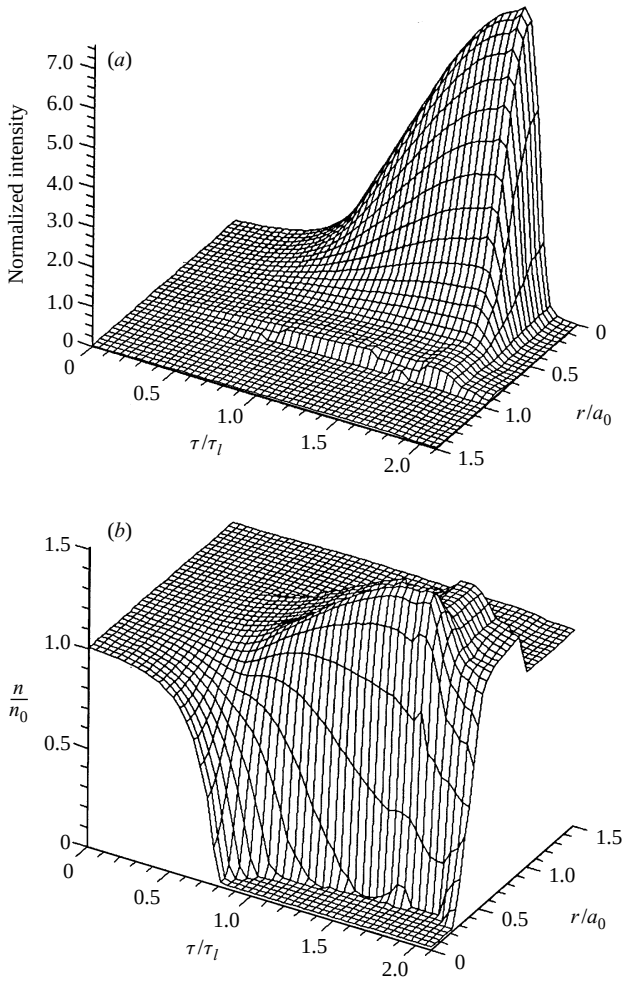


Figure 4. The intensity (a) and the corresponding electron-density profile (b) after the pulse has travelled a distance of L_R. The width and length of the channel are increased compared with Fig. 3.

self-focusing and self-guiding of the laser pulse. On all figures, for convenience, the origin of the τ/τ_l axis is displaced to the left by a distance of 1.8. Significant expulsion of electrons out of the axial region is clearly seen even at a distance of $0.4L_R$ (Fig. 2). The process of electron ejection is accompanied by the formation of an ion channel, and at a distance of $0.5L_R$ it is practically complete, as can be seen in Figs 3–5. Even at a distance of L_R, the length of the channel exceeds one-half of the initial length of the laser pulse, and the width is approximately one-third of its initial width in vacuum. The length of the leading edge of the pulse is nearly halved on propagation. Modulation of the envelope of the edge also occurs (Fig. 4). This modulation is connected with self-focusing and erosion of HF wave beam. At a distance of $8L_R$ (Fig. 5), the length of the laser pulse became nearly equal to the length of the ion channel.

With distance, the rear edge of the laser pulse is also shortened. But, owing to the abrupt cut-off of the rear edge, its shortening is less noticeable, so the

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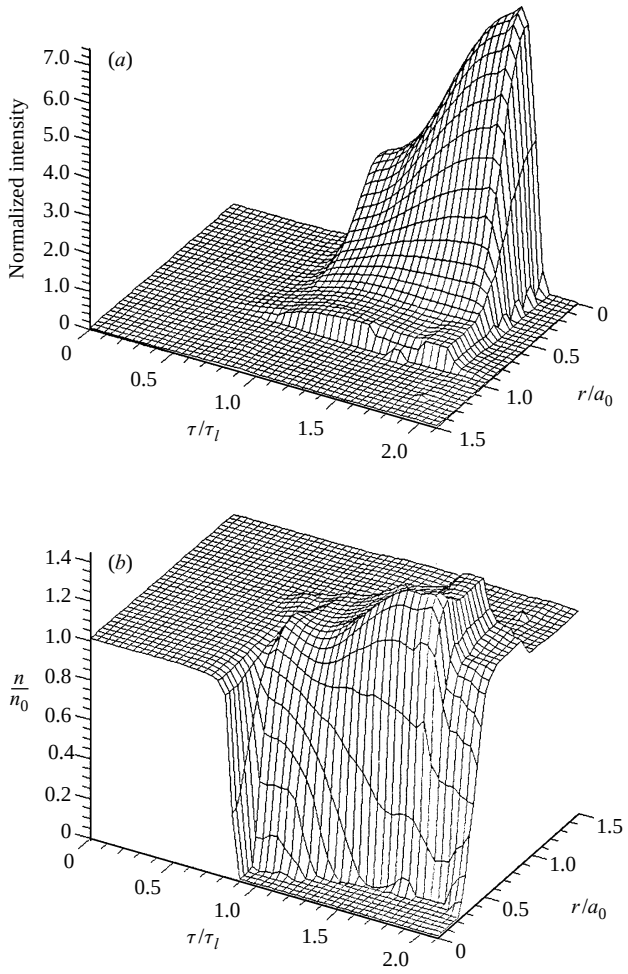


Figure 5. The intensity (a) and the corresponding electron-density profile (b) after the pulse has travelled a distance of $8L_R$. Strong redistribution of the pulse intensity and the electron density at the leading front of the pulse is evident.

width of the ion channel at the rear edge remains almost unchanged. Thus the time τ_r is insufficient for the electrons ejected by the leading edge to the periphery of the ion channel to be redistributed appreciably in the radial direction.

According to Fig. 6, a longitudinal electrical field E_z is generated at the leading edge of the pulse, and increases monotonically towards the end of the rear edge. This field is given in dimensionless units by

$$E_z = \frac{e}{mc\omega_0}\left(\frac{\omega_0}{\omega_p}\right)^2 \frac{1}{\omega_0 \tau_l} E_A. \tag{44}$$

The width of the accelerating field is increased towards the rear edge in the radial direction; at a distance of $\frac{1}{2}a_0$ it changes its polarity, and as $r \to \infty$ it becomes close to zero.

To obtain a qualitative estimate, we take the following characteristic values

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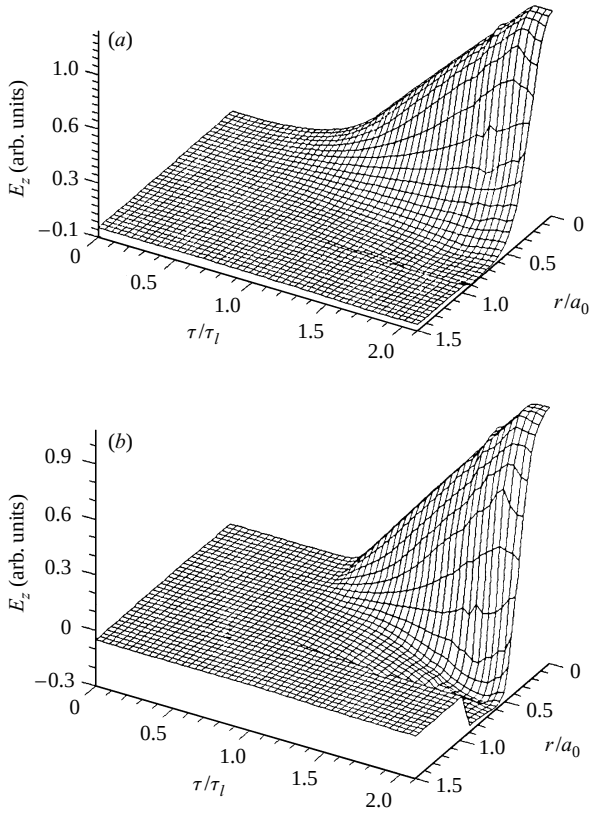


Figure 6. The axial electrostatic field generated inside the pulse after the pulse has travelled (*a*) a distance of $0.5L_R$ and (*b*) a distance of $8L_R$. The field is mainly concentrated in the channel, and tends to zero as $r \to \infty$.

of the parameters: $\lambda_0 \approx 1.06\,\mu$m, $\tau_l \approx 0.2$ ps and $\omega_p^2 \approx 3 \times 10^{27}$ s^{-2}. From Fig. 5, it can be seen that, before filamentation and damage, the laser pulse travels a distance of at least $8L_R$. For the above parameters, according to Fig. 6(*b*) and (44), we obtain $E_A \approx 10$ GV cm^{-1} for the accelerating field on the axis. An electron with the accelerating phase of such field gains an energy of $8L_R eE_A \approx 10$ GeV at a distance of $8L_R \approx 1$ cm.

Thus it follows from our considerations that at $0.4L_R$ there is clearly self-guiding of the laser pulse. Figure 5 shows the distribution of the laser-pulse intensity and the corresponding electron-density profile when the pulse is moved a distance $8L_R$ in the plasma. It can be seen that at this distance the self-guiding regime of propagation is maintained without appreciable change in the width of the ion channel, and the pulse, followed by the ion channel, continues to move with finite length and reduced width.

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