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Plasma-field structures during relativistic laser interaction with overdense plasmas at finite electron temperatures

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Abstract. A simple model taking into account the effect of electron temperature is derived to define the plasma-field structures which may arise during relativistically intense laser interaction with overdense plasmas. We show that there exist multilayer solutions with electron cavitation, allowing for both relativistic and ponderomotive nonlinearities. The influence of finite electron temperature on such structures is studied. Examples of these plasma-field structures for the cases of an infinite plasma and a plasma layer are presented.

PACS. 52.38.-r Laser-plasma interactions – 52.35.Mw Nonlinear phenomena

1 Introduction

Recent developments in laser technology have opened up the possibility of exploring laser-matter interactions in previously unattainable regimes [1]. Intensities on the order of 10^{21}–10^{22} W/cm^2 and higher can now be achieved, thus imparting strong momentum to theoretical investigations of phenomena occurring under such extreme conditions, when electron motion becomes ultrarelativistic and new regimes of interaction may appear [2], suggesting that goals such as compact sources for short-pulse coherent X-ray radiation [3–5], the fast ignitor concept for inertial confinement fusion (ICF) [6], and laser-plasma based particle accelerators [7–9] might soon be within reach. However, a major effort is still required both numerically and analytically in order to understand the nonlinear phenomena that arise in the presence of such extremely high intensities of electromagnetic radiation. Good analytical insight is also needed in order to make numerical simulations possible and to interpret their results.

It is well known that at laser intensities exceeding 10^{18} W/cm^2 electron motion in the laser field becomes relativistic, which leads to significant changes in the dynamics of interaction of such radiation with plasmas. In particular, owing to the Lorentz mass increase the effective plasma frequency $\omega_p = (4\pi e^2 N_e/m_e)^{1/2}$ (N_e, e, m_e are the electron density, charge and mass, respectively) decreases and the plasma that was initially nontransparent at this frequency may become transparent [10]. In addition, the role of the magnetic component of the Lorentz force acting on the electrons increases considerably in the relativistic regime because this force is proportional to the ratio v_e/c, where v_e is electron velocity. This force also gives rise to ponderomotive action of radiation on plasma which, at such high intensities, plays a decisive role in plasma dynamics.

As was mentioned above, the interaction of superintense laser radiation with plasma is a very complicated problem and the main instrument for theoretical research in this field is numerical simulation using very powerful computers. However, interpretation of the results of these studies and formulation of the problem demand physical intuition based on simple analytical approaches. Therefore, of particular significance in terms of understanding the physics of the essentially nonlinear processes are the results obtained analytically in some model problems.

In the present work we consider one-dimensional quasi-stationary plasma-field structures that can be generated in overdense plasmas normally irradiated by a relativistically strong circularly polarized laser pulse. This research is a logical continuation of the work done in [11–16], in which analogous studies were carried out in the framework of the so-called "cold" plasma approximation, i.e. when the energy of electrons' thermal motion is negligible compared to the energy of their oscillatory motion in the laser field, we can assume the electron temperature to be zero. Within this approximation even relatively complex plasma-field structures can be found, in particular, localized multilayer structures with electron cavitation in infinite plasmas [14] and/or Bragg-like cavitated electron structures in semi-infinite plasmas describing the penetration of laser radiation into overdense plasmas to a finite depth only [15]. However, these solutions were found in a class of discontinuous functions, which means that from the physical point of view their potential realizability should be additionally proven. This can be done in two different ways: the

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first is by employing non-stationary equations and solving, in fact, a more complex time-dependent problem, for example, as in [17–19]. The second is by using some extension but within the framework of a quasi-stationary model; in particular, by introducing finite plasma temperature we are able to overcome any discontinuities in density distributions. Of course it cannot provide us with a dynamical scenario of laser-plasma interaction, but this way is much simpler compared to the first one and gives us not only useful information on realizable plasma-field structures but also provides their dependence on electron temperatures, which can be fairly large as well. It should also be noted that from the theoretical point of view, in the limit of an electron temperature of roughly zero, a small parameter in the higher derivative is neglected. In the current work we introduce this small parameter characterizing electron temperature that provides us with a unique way of finding plasma-field structures uniformly. In particular, it is shown that not all plasma-field structures previously found at $T_e = 0$ can occur. For example, anti-symmetrical structures with respect to the field distribution [14] are now forbidden. It is also useful to note that in a linearly polarized wave, inherent $[\mathbf{j}, \mathbf{B}]$ heating at relativistic intensities (see, e.g., [20]) complicates the interaction dynamics and makes them significantly different from the case of circular polarization. In the latter, electron heating can be significantly suppressed [21], thus allowing to consider the temperature of electrons as an parameter in the theoretical model of relativistic laser-plasma interaction.

The paper is organized as follows. In Section 2, the basic model is described and some limiting cases are discussed. Section 3, which consists of two subsections, is devoted to the plasma-field structures. In the first subsection, solutions for the case of an infinite homogeneous plasma are presented, whereas in the second subsection we study the case of incident laser radiation on a plasma layer where the field and electron density distributions are defined. These results illustrate the interplay between the different effects and are in qualitative agreement with the results obtained previously, for example, in [14,15]. Finally, in Section 4 the conclusions are summarized.

2 Basic model

Let us consider the interaction between an overdense plasma and an ultrahigh power laser with comparatively long pulses so that the pulse width exceeds the electron oscillation time ($\tau \gg \omega_{pe}^{-1}$) but is less than the ion response time ($\tau \ll \omega_{pi}^{-1}$). This makes it possible to treat the ions as a stationary neutralizing background. Given the assumptions about the characteristic time scales of the interacting pulse and the plasma, we expect that stationary plasma-field structures can emerge, since the electron fluid should have time to approach a quasisteady state. The question of whether these structures actually appear on a time scale that effectively allows us to disregard other effects and complications was in fact confirmed by numerical simulations in [17,18]. In order to arrive at the quasi-steady model of relativistic laser plasma interactions, we reduce the system of Maxwell's equations and the relativistic hydrodynamic equations describing the motion of an electron fluid by adopting the Coulomb gauge $\mathrm{div}\mathbf{A} = 0$ and representing the vector potential in the following dimensionless form

$$\frac{e\mathbf{A}(z,t)}{m_0c^2} = a(z)\mathrm{Re}\left[e^{i\vartheta(z)}\left(\mathbf{x_0} + i\mathbf{y_0}\right)e^{i\omega t}\right], \qquad (1)$$

where $\mathbf{x_0}, \mathbf{y_0}$ are unit vectors along the x- and y-axes, z is the propagation coordinate, ω is the laser frequency, and m_0 is the electron rest mass. As was shown in [12,13], Maxwell's equations in plasma in this case may be reduced to the following equations

$$\frac{d^2a}{d\zeta^2} - \frac{s^2}{a^3} + \left(1 - \frac{n_0 n}{\gamma}\right)a = 0, \qquad (2)$$

$$\frac{d^2\phi}{d\zeta^2} = n_0(n - Z_i n_i), \qquad (3)$$

where $\zeta = \omega z/c$; $\phi = e\varphi/m_0c^2$ is the normalized scalar potential; the normalized electromagnetic energy flux density is $s = -a^2(d\vartheta/d\zeta) = \text{const.}$ that is the parameter of the problem; $n_0 = 4\pi e^2 N_0/m_0\omega^2$ is the plasma overdensity parameter equal to the ratio of the initial electron density N_0 to its critical value at a given laser frequency; Z_i is the ion charge (all the ions are taken to have the same charge); $n = N_e/N_0, n_i = N_i/N_0$ are the normalized electron and ion densities, respectively; and $\gamma = \sqrt{1+a^2}$ is the electrons' relativistic factor (i.e., the ratio of their energy to their rest energy; thermal energy is neglected here). Thus, the above simplification has enabled us to reduce Maxwell's equations to a nonlinear Helmholtz equation for the vector potential's complex amplitude (2) and Poisson's equation for the scalar potential (3). The first equation describes the transverse field of the incident laser wave, and the second one describes the longitudinal electrostatic field in the plasma caused by charge separation.

The relativistic hydrodynamic equations for a plasma with uniform bulk electron temperature T_e and immobile ions can be written as

$$\frac{\partial N_e}{\partial t} + \mathrm{div}(N_e\mathbf{v}) = 0, \qquad (4)$$

$$\frac{\partial \mathbf{p}}{\partial t} + (\mathbf{v}\nabla)\mathbf{p} = -T_e\nabla N_e + \frac{e}{c}\frac{\partial \mathbf{A}}{\partial t} + e\nabla\varphi - \frac{e}{c}[\mathbf{v}\cdot\mathrm{rot}\mathbf{A}], \qquad (5)$$

$$\mathbf{p} = \gamma m_0 \mathbf{v}. \qquad (6)$$

Here $\mathbf{v}, \mathbf{p}$ are the mean velocity and momentum of the electron component, respectively. This set of equations may also be reduced in a one-dimensional geometry. As a result, in the quasisteady case it reduces to two equations, the first of which is the conservation law of a canonical electron momentum in the transverse direction:

$$\mathbf{p}_\perp - \frac{e}{c}\mathbf{A} = 0. \qquad (7)$$

The second equation is a simple balance of forces acting on electrons:

$$\frac{d\phi}{d\zeta} - \frac{d\gamma}{d\zeta} - \mu\frac{1}{n}\frac{dn}{d\zeta} = 0. \quad (8)$$

The first term here corresponds to the electrostatic force of charge separation in plasma, the second one to the ponderomotive force of the laser field, and the third term corresponds to the internal pressure forces arising due to the temperature effect, where the parameter $\mu = T_e/m_0c^2$ characterizes electron temperature. The continuity equation (4) is fulfilled automatically in this case, thus ensuring the plasma quasineutrality condition.

3 Steady state plasma-field structures

Next, we consider the two most interesting cases which may arise in relativistic laser interactions with overdense plasmas. The first one is the case of an infinite homogeneous plasma that is of interest firstly from the fundamental point of view of the basic plasma-field structures which could be generated during the laser-plasma interactions. Secondly, it could be important in terms of practical applications since it could provide an insight into the physics of laser-foil interactions and their use for the ion acceleration problem.

3.1 Infinite plasma

We now consider solutions of the set of equations (2), (3) and (8) in the case of infinite homogeneous plasma. Assuming $Z_i n_i = 1$, it is convenient to integrate equation (8) and express the electron density as follows

$$n = \exp\frac{\phi - \gamma + 1}{\mu}. \quad (9)$$

Implying that we are interested in localized solutions only, for which $a \to 0, \phi \to 0, n \to 1$ is fulfilled for $\zeta \to \pm\infty$. In addition, $s = 0$ also holds for localized solutions (nonlinear standing wave structures) in an infinite plasma (see also in [11,14]). Substituting (9) into the two remaining equations yields the following fourth-order system of ordinary differential equations

$$\frac{d^2a}{d\zeta^2} = \left(\frac{n_0}{\sqrt{1+a^2}}\exp\frac{\phi - \sqrt{1+a^2}+1}{\mu} - 1\right)a, \quad (10)$$

$$\frac{d^2\phi}{d\zeta^2} = n_0\left(\exp\frac{\phi - \sqrt{1+a^2}+1}{\mu} - 1\right). \quad (11)$$

Note that solutions of equations (10) and (11) can easily be found numerically since they have the following integral

$$\frac{1}{2}\left(\frac{da}{d\zeta}\right)^2 + \frac{1}{2}a^2 - \frac{1}{2}\left(\frac{d\phi}{d\zeta}\right)^2 - n_0\phi - \mu n_0 \exp\frac{\phi - \sqrt{1+a^2}+1}{\mu} = \text{const}. \quad (12)$$

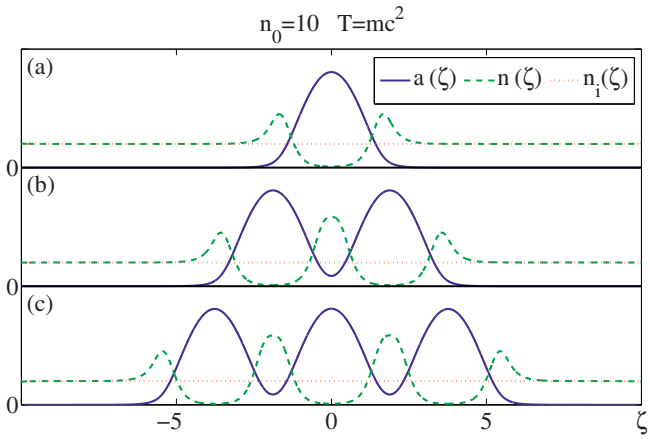

Fig. 1. (Color online) Examples of localized solutions of the set of equations (10), (11). (a) First mode with one field antinode. (b) Second mode, two field antinodes. (c) Third mode, three field antinodes.

This integral simplifies the search for localized solutions to the problem of interest. Indeed, let us set the problem of finding localized solutions which are symmetric with respect to the point on the $\zeta = 0$ plane. The following relation

$$\left.\frac{da}{d\zeta}\right|_{\zeta=0} = 0 \quad \left.\frac{d\phi}{d\zeta}\right|_{\zeta=0} = 0$$

is fulfilled for such solutions. Then for the integral (12) we obtain

$$\frac{1}{2}a^2(0) - n_0\phi(0) - \mu n_0 \times \left(\exp\frac{\phi(0) - \sqrt{1+a^2(0)}+1}{\mu} - 1\right) = 0.$$

Thus using the so-called shooting method for solving ordinary differential equation, for which we now have one fitting parameter (for example, $\phi(0)$) we can find a localized solution of the set (10), (11). Our studies show that in such a formulation there may exist a family of solutions differing by the number of field antinodes. Examples for the first three modes are given for $\mu = 1$ in Figure 1. In addition, the integral (12) allows us to show that there are no antisymmetric solutions within the framework of the set of equations (10), (11) for which the relation

$$a\Big|_{\zeta=0} = 0 \quad \left.\frac{d\phi}{d\zeta}\right|_{\zeta=0} = 0$$

is fulfilled in the $\zeta = 0$ plane. At the same time, such solutions in a cold plasma (see, e.g., [16]) can be found. This is not surprising either physically or mathematically. Indeed, in terms of the mathematics, we have increased the order of the studied system of equations, which may lead to appreciable changes in solution topology and phase plane in general. In terms of physics, introduction of nonzero temperature results in appearance of electrons that can move from one field antinode to another during thermal

motion, thus providing their interaction. This interaction most likely leads to the non-stationarity of antisymmetric distributions.

Solutions constructed in an infinite plasma point to some principal features of the regime of interaction of a superintense laser pulse with overdense plasma. These are appreciable redistribution of electron density under the action of ponderomotive force and the related longitudinal electrostatic field being approximately equal to the laser field, as well as existence of structures with several electron layers separated by gaps almost devoid of electrons (Figs. 1b, 1c).

3.2 Plasma layer

Let us now consider the case of a bounded plasma layer that is more important for practical reasons. We will restrict our consideration to the one-sided irradiation of the plasma. Then, it is convenient to construct solutions starting at some point outside the plasma layer, from the unirradiated side rather far from the plasma so that the conditions $n, \phi', a' \ll 1$ (the primes denote differentiation along the coordinate) can be regarded to be met at this point. Note that $a \approx \sqrt{s}$, and the function $\phi(\zeta)$ is generally unbounded, since in the one-dimensional geometry the potential is known to behave logarithmically at large distances. By introducing a new variable e that denotes a longitudinal electrostatic field and is defined by

$$e(\zeta) = -\frac{d\phi}{d\zeta}$$

and eliminating from the basic equations (2), (3) and (8) the variable ϕ for which boundary conditions are not clear, we can obtain the following set

$$\frac{d^2a}{d\zeta^2} = \frac{s^2}{a^3} - \left(1 - \frac{n_0 n}{\sqrt{1+a^2}}\right) a \tag{13}$$

$$\frac{de}{d\zeta} = n_0 \left(Z_i n_i(\zeta) - n\right) \tag{14}$$

$$\mu \frac{dn}{d\zeta} = -n \left(e + \frac{d\sqrt{1+a^2}}{d\zeta}\right) \tag{15}$$

here $n_i(\zeta)$ is a given function specifying the initial plasma distribution in space. We will further restrict our consideration to the case of a plane homogeneous layer of thickness L so that this function will have the form

$$n_i(\zeta) = \begin{cases} 0 & \text{for} \quad \zeta < 0,\ \zeta > \omega L/c, \\ 1/Z_i & \text{for} \quad 0 < \zeta < \omega L/c. \end{cases} \tag{16}$$

By analogy with (10), (11) the system (13)–(15) has an integral analogous to (12) which, in the case under consideration, has the following form

$$\frac{1}{2}\left(\frac{da}{d\zeta}\right)^2 + \frac{1}{2}a^2 - \frac{1}{2}\left(\frac{s}{a}\right)^2 - \frac{1}{2}e^2 - \mu n_0 n - n_0 Z_i \int e n_i(\zeta) d\zeta = \text{const}. \tag{17}$$

We see that for $n_i(\zeta) \neq 0$ the integral (17) contains integration in explicit form, which restricts its applicability. However, it may be used in the region fully devoid of ions, where the last term in the left-hand side of the expression (17) is identically equal to zero.

As was mentioned above, we will construct solutions by direct numerical integration starting at some point ζ_{st} that is rather far from the plasma layer on its unirradiated side. Suppose that $a = \sqrt{s} + \alpha$ at this point and the inequalities $\alpha(\zeta_{st}) \ll \sqrt{s}$, $n(\zeta_{st}) \ll 1$, $e(\zeta_{st}) \ll 1$ are fulfilled. In addition, $n_i(\zeta_{st}) = 0$ as this point is outside the plasma layer. Then in a standard manner, in the neighborhood of the point ζ_{st} making an expansion of equations (13)–(15) and keeping all terms first order in the small quantity of $\alpha/\sqrt{s}$ we can arrive at

$$\frac{d^2\alpha}{d\zeta^2} = -4\alpha + n_0 n \sqrt{\frac{s}{1+s}} \tag{18}$$

$$\frac{de}{d\zeta} = -n_0 n \tag{19}$$

$$\mu \frac{dn}{d\zeta} = -n\left(e + \sqrt{\frac{s}{1+s}}\frac{d\alpha}{d\zeta}\right). \tag{20}$$

As the system (18)–(20) is still nonlinear, its solution is rather complicated, but assuming that the inequality

$$e \gg \sqrt{\frac{s}{1+s}}\frac{d\alpha}{d\zeta} \tag{21}$$

is satisfied, equations (19) and (20) can be separated from equation (18), which allows one to find their solutions. Indeed, if we divide the second of these equations by n, differentiate it and substitute $de/d\zeta$ from the first equation, we obtain one second-order equation for n:

$$\mu \frac{d^2 n}{d\zeta^2} - \mu \frac{1}{n}\left(\frac{dn}{d\zeta}\right)^2 - n_0 n^2 = 0. \tag{22}$$

Let us seek a solution to this equation of the form

$$n = C(\zeta - \zeta_0)^\beta$$

and obtain

$$\mu C\beta\left(\beta - 1\right)(\zeta - \zeta_0)^{\beta-2} - \mu C\beta^2(\zeta - \zeta_0)^{\beta-2} - n_0 C^2 (\zeta - \zeta_0)^{2\beta} = 0.$$

This equality must be identically fulfilled for all values of ζ; therefore, the conditions

$$\beta = -2, \qquad C = \frac{2\mu}{n_0}$$

must be fulfilled. Finally the solution of equation (22) takes on the form

$$n(\zeta) = \frac{2\mu}{n_0(\zeta - \zeta_0)^2}. \tag{23}$$

From equation (19) we find

$$e(\zeta) = \frac{2\mu}{\zeta - \zeta_0}, \tag{24}$$

and from equation (18) follows

$$\alpha(\zeta) = \frac{\mu}{2(\zeta - \zeta_0)}\sqrt{\frac{s}{1+s}}. \tag{25}$$

We can now readily verify that the inequality (21) that was supposed to be met does hold sufficiently far from the point $\zeta = \zeta_0$. Note that all the integration constants in the derivation of equations (23)–(25) were chosen so that the quantities α, e, n vanish at infinity.

Generally speaking, in the construction of the solution ζ_0 is an unknown quantity that determines a whole family of solutions. There arises the question of which of these solutions is correct. The answer to this question is quite trivial. One must choose the solution for which the conditions of overall plasma quasi-neutrality are fulfilled:

$$\int_{-\infty}^{+\infty} (Z_i n_i(\zeta) - n(\zeta))\, d\zeta = 0. \tag{26}$$

Actually, the problem of constructing a solution reduces to sorting all solutions by parameter ζ_0 and finding the one that satisfies condition (26). Moreover, there is no need to introduce ζ_0 in this case. Indeed, from (23)–(25) we can find the following relations

$$n = \frac{e^2}{2\mu n_0} \tag{27}$$

$$\alpha = \frac{e^2}{8\mu}\sqrt{\frac{s}{1+s}} \tag{28}$$

$$\frac{d\alpha}{d\zeta} = -\frac{e^3}{8\mu^2}\sqrt{\frac{s}{1+s}}. \tag{29}$$

Thus, we have excluded the parameter ζ_0.

Now the solution procedure will be the following. We choose a certain point $\zeta_{st} > \omega L/c$ at which we set the value of variable e. The values of a, $da/d\zeta$ and n at this point are found from the relations (27)–(29) taking into account that $a = \sqrt{s}+\alpha$. Thus, we have boundary conditions for the system (13)–(15) at the point ζ_{st}. We solve the system of equations and check whether the equality (26) is fulfilled. If it is fulfilled, then it is the required solution; if the inequality is not fulfilled, we change the value of the variable e at the point ζ_{st} and repeat the procedure. We proceed in this fashion until we find the right solution. In view of the complexity of the system its solution was done numerically by manually sorting the value of $e(\zeta_{st})$. (There are some ideas of how this sorting might be automated, but they have not been implemented for lack of necessity.)

Now a few words about the parameter s. As we have mentioned above, it represents the electromagnetic energy flux density and is equal to the intensity of radiation passing through the layer. In our problem it determines the amplitude of the wave incident on the layer. Thus, if we have to construct a function of the incident wave amplitude, then we must sort parameter s for each of its values using the procedure described above.

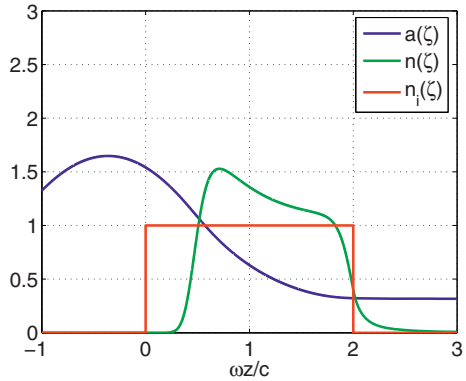

Fig. 2. (Color online) Example of solution of system (13)–(15) in a plasma layer of finite thickness. Parameters of the problem: $\mu = 0.01$, $n_0 = 2$, $\omega L/c = 2$, $s = 0.1$. The corresponding amplitude of incident radiation $a_i = 0.85$.

An example of a solution constructed using the described algorithm is presented in Figure 2. The parameters of the problem were taken to be $n_0 = 2$, $s = 0.1$, $\omega L/c = 2$. As soon as the solution has been constructed one can find the amplitude of the incident pulse which, in a real situation, is a true parameter of the problem specified in the course of an experiment. For this the values of a and $da/d\zeta$ must be known at an arbitrary point in the vacuum on the irradiated side of the plasma layer. Indeed, at this point there are incident and reflected waves, hence we can write the following equality

$$a(\zeta)e^{i\vartheta(\zeta)} = \hat{a}_i e^{-i\zeta} + \hat{a}_r e^{i\zeta}, \tag{30}$$

where $\vartheta(\zeta)$ was introduced in equation (1) and $\hat{a}_i$ and $\hat{a}_r$ denote complex amplitudes of the incident and reflected waves, respectively. By differentiating equation (30) and omitting from the resulting system the quantity $\hat{a}_r$ we obtain

$$\left[\frac{da}{d\zeta} + ia\left(\frac{d\vartheta}{d\zeta} - 1\right)\right] e^{i\vartheta} = -2\hat{a}_i e^{-i\zeta}. \tag{31}$$

Taking the modulus of this expression and bearing in mind that parameter s was introduced in the form (see [16])

$$s = -a^2\frac{d\vartheta}{d\zeta}$$

we finally obtain an expression for the real amplitude of the wave incident on the layer:

$$a_i = \frac{1}{2}\left[\left(\frac{da}{d\zeta}\right)^2 + a^2\left(\frac{s}{a^2} + 1\right)^2\right]^{1/2}. \tag{32}$$

It is also interesting to compare solutions with different values of electron temperature μ. If parameter s is fixed, then the value of the incident wave amplitude will be different for different values of μ. Therefore, such a comparison will not be physically correct. For a more correct comparison we fitted the parameter s to each value of μ so as to ensure equal incident wave amplitudes in all the cases.

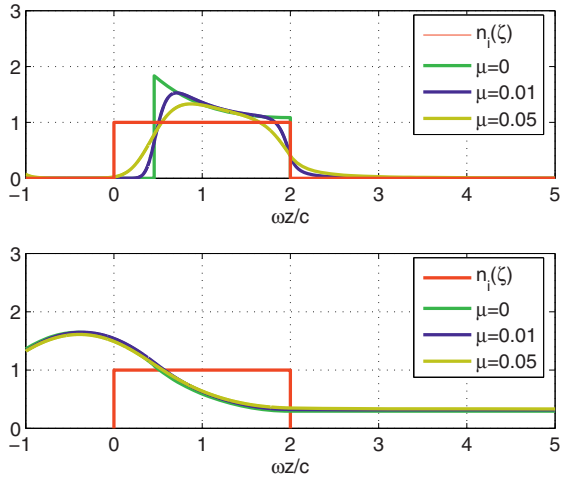

Fig. 3. (Color online) Comparison of solutions of system (13)–(15) for different values of μ but equal incident wave amplitudes. Distributions of electron concentration are given on top, and of laser field at the bottom. The solution for $\mu = 0$ was constructed by the method described in [16].

The result is illustrated in Figure 3. Note that the solutions obtained for $\mu = 0$ by the method described in [16] are also shown in this figure. The characteristic feature of this solution is enhancement of electron layer border blurring with increasing temperature, which is a natural and expected effect. In general, it should be noted that solutions of the system (13)–(15) with $0 < \mu < 0.01$ differ from the solutions with $\mu = 0$ obtained in the work [16] only slightly and are comparable in terms of complexity of their construction. This enables us to conclude that both the approaches to finding steady-state distributions are equivalent. The only drawback of the solutions with nonzero temperature is divergence of the value of the scalar potential at infinity. However, this divergence is logarithmic and may be eliminated manually by truncating solutions at some distance from the plasma layer.

In [16] it was shown that the so-called resonator-like structures in which electron layers were separated merely by ion gaps playing the role of walls may exist in a layer of finite thickness. Generally speaking, structures with an arbitrary number of electron layers separated by ion gaps may exist in a plasma layer. However, a special algorithm is needed for constructing solutions with N electron layers as such solutions are discontinuous, with the number of solution discontinuities being proportional to N. This circumstance strongly complicates construction of such solutions, whereas under the action of temperature the solutions become smooth and all the family may be obtained identically by simply varying the values of s and $e(\zeta_{st})$. An example of a solution with three electron layers, one of which is outside the plasma layer, is given in Figure 4.

4 Conclusion

To conclude, we have analysed quasisteady plasma-field structures that can be generated during relativistically intense laser interactions with overdense plasmas. By introducing the electron temperature effect we show that a uniform method can now be suggested for finding these plasma-field structures. This method was applied to the two most interesting cases of an infinite homogeneous plasma and a plasma layer where the corresponding structures were found. The presented solutions may also be used for analysis of possible ways of producing strong longitudinal fields in solid-state targets for the problem of ion acceleration.

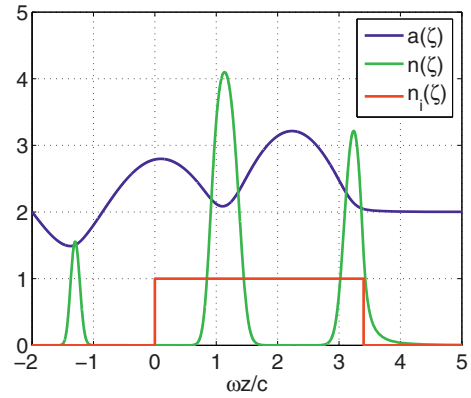

Fig. 4. (Color online) Example of solution of system (13)–(15) with three electron layers. Layer parameters: $\mu = 0.01$, $n_0 = 2$, $\omega L/c = 3.4$; incident wave amplitude $a_i = 2.1$.

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