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Instability of a plasma produced by laser-induced ionisation of a target

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***Abstract*.** **The development of a longitudinal – transverse instability in a plasma with the anisotropic velocity distribution of electrons produced by the tunnelling ionisation of a target by a short laser pulse is studied. The dependences of the growth rate of this longitudinal – transverse instability on the wave number and the propagation angle of perturbations are investigated. It is shown that the increasing longitudinal electrostatic field in a rather broad angular region is comparable with the exciting magnetic field. It is pointed out that the longitudinal – transverse instability of the anisotropic plasma can lead to the additional absorption of laser radiation by the nonthermal fluctuations of the electron density.**

Keywords: *plasma instability, tunnelling ionisation, short laser pulse.*

1. Introduction

Subpicosecond light pulses generated by modern lasers can ionise materials due to the tunnelling effect [1]. The properties of tunnelling ionisation and a plasma produced due to ionisation were investigated both experimentally [2 – 4] and theoretically [5 – 9]. It is well known that a plasma produced due to tunnelling ionisation of atoms by laser radiation has the anisotropic energy distribution of electrons. Thus, for example, the energy distribution of electrons ejected upon the ionisation of atoms by linearly polarised light proves to be similar to the two-Maxwell distribution, while in the case of circularly polarised radiation, this distribution is characterised by a change in the projection of the electron velocity on the polarisation plane of radiation, which considerably exceeds the characteristic velocity along the laser beam.

The anisotropy of the velocity distribution of electrons in a photoionised nonequilibrium plasma leads to the development of instabilities, which considerably affect the plasma properties and absorption of laser radiation by plasma. Such instabilities are observed for laser radiation both in the optical and X-ray ranges, although ionisation in the first case is caused by tunnelling, while in the second case, it occurs due to the usual photoeffect. In the case of anisotropy caused by tunnelling ionisation, the main attention was paid so far to the Weibel instability [10], similarly to the classical instability in the plasma with the anisotropic temperature distribution [11]. This instability develops across the direction along which temperature is maximal and is characterised by the increase in the quasi-stationary transverse electromagnetic field.

The generation of the magnetic field $\boldsymbol{B}$ in the plasma can directly affect the depolarisation of laser radiation, whereas the absorption of this radiation can be changed indirectly due variations in the transfer of absorbed energy or due to density perturbations caused by nonlinear effects (proportional to B^2). At the same time, it is known that in the general case of the propagation of instable perturbations at an arbitrary angle to the anisotropy axis, the instability of the plasma with the anisotropic temperature [11] is longitudinal – transverse and, along with a magnetic field, causes the increase in the longitudinal electric field as well, or, which is the same, in perturbations of the electron density. This also concerns, of course, the anisotropic energy distribution function of electrons caused by the field ionisation of the medium in the tunnelling regime. It is this fact that we pay attention to in our paper, pointing to the possibility of the development of density perturbations due to the instability of the plasma with the anisotropic energy distribution function of electrons in the region of laser radiation cut-off, which can result in the increase in the light pulse absorption due to such perturbations. Note that the increase in the absorption of laser radiation caused by small-scale density perturbations was discussed in papers [12 – 15].

We investigated the properties of the longitudinal – transverse instability of an anisotropic plasma produced due to tunnelling ionisation by a short laser pulse in the optical region. For definiteness, we consider circularly polarised light, which admits a less cumbersome description of the instability compared to the case of linearly polarised light. We also assume that the energy of electrons released in ionisation is high enough for neglecting electron – ion collisions. The angular dependence of the instability growth rate and the relation between the amplitudes of the longitudinal and transverse components of the excited electromagnetic field were investigated. We also discussed the possibility of the instability development from the initial level of density perturbations, which considerably exceeds the spontaneous level and is caused by the formation of a standing wave in the reflection region.

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2. Solution of the dispersion equation

The energy distribution function of the electrons produced due to the tunnel ionisation of materials by short laser pulses is strongly anisotropic [5]. The anisotropy is preserved until the relaxation and redistribution of particles over velocities occur due to their nonlinear interaction and (or) collisions. In this case, the problem of the stability of the plasma produced due to photoionisation by a short laser pulse can be formulated. This problem can be most simply solved in the case short-wavelength perturbations for which the characteristic wavelength is small compared to the characteristic scale of plasma inhomogeneities (the focal spot size and the gradient length) and, therefore, the plasma can be treated as homogeneous. Below, we assume that this condition is fulfilled.

For definiteness, we consider the case of circularly polarised laser radiation, when the direction of the plasma anisotropy coincides with the laser pulse propagation direction (hereafter, denoted by the unit vector $\boldsymbol{n}$) along which the characteristic velocity $v_{\parallel}$ of electrons is considerably smaller than the characteristic velocity $v_{\perp}$ in the transverse plane (in the plane of the polarisation vector of laser radiation). By neglecting a small scatter of velocities $\boldsymbol{v}$, we can represent the velocity distribution function of electrons in the approximate form [5]

$$f_{\mathrm{e}}(\boldsymbol{v}) = \frac{n_{\mathrm{e}}}{\pi}\,\delta(v_{\parallel})\,\delta(v_{\perp}^2 - v_0^2). \quad (1)$$

Here, $v_0 = eE_0/m\omega_0$ is the rate of electron oscillations in the laser wave field; e and m are the electron charge and mass, respectively; E_0 and ω_0 are the laser-field amplitude and frequency, respectively; and n_{e} is the electron concentration in the plasma. It is also assumed that the electric field strength is high enough to produce tunnelling ionisation, but, however, is insufficient to induce relativistic effects, i.e. $v_0^2 \ll c^2$, where c is the speed of light.

Consider the standard dispersion equation for the plasma consisting of electrons with velocity distribution (1) and ions, which we treat as an immobile neutralising background [11]:

$$\Lambda(\omega, \boldsymbol{k}) \equiv \det|k^2\delta_{ij} - k_i k_j - (\omega^2/c^2)\epsilon_{ij}^{\mathrm{e}}(\omega, \boldsymbol{k})| = 0, \quad (2)$$

where $\boldsymbol{k}$ and ω are the wave vector and the perturbation frequency of the electron density;

$$\epsilon_{ij}^{\mathrm{e}}(\omega, \boldsymbol{k}) = \delta_{ij} - \frac{\omega_{\mathrm{pe}}^2}{\omega^2}\,\delta\epsilon_{ij}(\omega, \boldsymbol{k})$$

is the electron component of the plasma permittivity tensor expressed in terms of the electron velocity distribution function (1)

$$\delta\epsilon_{ij}(\omega, \boldsymbol{k}) = \delta_{ij} + \frac{1}{n_{\mathrm{e}}}\int \mathrm{d}^3 v f_{\mathrm{e}}(\boldsymbol{v})\,\frac{(k_i v_j + k_j v_i)(\omega - \boldsymbol{k}\boldsymbol{v}) + k^2 v_i v_j}{(\omega - \boldsymbol{k}\boldsymbol{v})^2};$$

and $\omega_{\mathrm{pe}} = (4\pi e^2 n_{\mathrm{e}}/m)^{1/2}$ is the electron plasma frequency.

Consider the possibility of the instability growth caused by aperiodic perturbations $\omega = \mathrm{i}\gamma$ (γ is the instability growth rate). Let us direct the wave vector $\boldsymbol{k}$ of these perturbations along the z axis, and direct the x and y axes in such a way that the vector $\boldsymbol{n}$ will lie in the xz plane. We denote the angle between this vector and the z axis by θ, assuming for definiteness that this angle changes in the interval $[0, \pi/2]$ ($\theta \leftrightarrow \pi - \theta$ due to symmetry). For this geometry of the problem, the components of the tensor $\delta\epsilon_{ij}$ take the form

$$\delta\epsilon_{ij} = \begin{pmatrix} \delta\epsilon_{xx} & 0 & \delta\epsilon_{xz} \\ 0 & \delta\epsilon_{yy} & 0 \\ \delta\epsilon_{zx} & 0 & \delta\epsilon_{zz} \end{pmatrix},$$

$$\delta\epsilon_{xx} = 1 + \cot^2\theta\left[1 - \beta\frac{\beta^2 + 2\sin^2\theta}{(\beta^2 + \sin^2\theta)^{3/2}}\right],$$

$$\delta\epsilon_{yy} = 1 + \frac{1}{\sin^2\theta}\left[1 - \frac{\beta}{(\beta^2 + \sin^2\theta)^{1/2}}\right], \quad (3)$$

$$\delta\epsilon_{zz} = \frac{\beta^3}{(\beta^2 + \sin^2\theta)^{3/2}},$$

$$\delta\epsilon_{xz} = \delta\epsilon_{zx} = \frac{\beta\cos\theta\,\sin\theta}{(\beta^2 + \sin^2\theta)^{3/2}},$$

where $\beta = \gamma/kv_0$. Equation (2) splits into two independent equations

$$k^2 - \frac{\omega^2}{c^2}\,\epsilon_{yy} = 0, \quad (4)$$

$$\left(k^2 - \frac{\omega^2}{c^2}\,\epsilon_{xx}\right)\epsilon_{zz} + \frac{\omega^2}{c^2}\,\epsilon_{xz}^2 = 0, \quad (5)$$

as in the case of the standard uniaxial anisotropy of the plasma [11].

Equation (4) describes the propagation of a transverse high-frequency electromagnetic wave. It can be easily shown that such waves are stable for any propagation angle θ, and we will not consider them in our study. In the case of an arbitrary propagation angle of perturbations, equation (5) describes longitudinal – transverse aperiodic oscillations, which can grow, i.e. the equation can have real positive solutions $\gamma(k, \theta) > 0$.

Dispersion equation (2) was solved in [5] only for the case $\theta = 0$, i.e. the propagation of perturbations was considered only along the laser beam. In this case, equation (5) is strongly simplified and gives the known expression for the growth rate [5]:

$$\gamma(k, 0) = \frac{kv_0\omega_{\mathrm{pe}}}{[2(\omega_{\mathrm{pe}}^2 + k^2c^2)]^{1/2}}, \quad (6)$$

which describes excitation of the non-potential Weibel instability corresponding to the appearance of a quasi-stationary magnetic field. When the angle θ is close to $\pi/2$, dispersion equation (5) is also simplified and its approximate solution, which gives the small value of $\gamma(k, \theta)$ compared to (6), has the analytic form

$$\gamma(k, \theta) \simeq \frac{kv_0\omega_{\mathrm{pe}}\cos^2\theta}{\omega_{\mathrm{pe}}^2 + k^2c^2}.$$

In the case of an arbitrary angle, the analytic solution for the instability growth rate is absent, and dispersion equation (5) was solved numerically. Below, for definiteness, the solution $\gamma(k, \theta)$ is illustrated for $v_0/c = 0.3$.

Figure 1a shows the dependence of the growth rate γ on the wave number for different angles between the direction of the anisotropy axis $\boldsymbol{n}$ and the wave vector $\boldsymbol{k}$. One can see that for $\theta = 0$ the instability has no the short-wavelength cut-off and its growth rate saturates for $k \to \infty$, which corresponds to (6). For nonzero angles θ of perturbation propagation, the dependence of the growth rate on the wave number changes. The maximum γ^* of the growth rate dependence on k appears, and there exists an 'optimal' wave number $k^*(\theta)$ corresponding to the most intense perturbation growth. Because the growth rate decreases with increasing the wave number, the characteristic short-wavelength 'cut-off' for k exists. As the cut-off value, we can use the wave number value at which the growth rate decreases compared to the maximal value by e times.

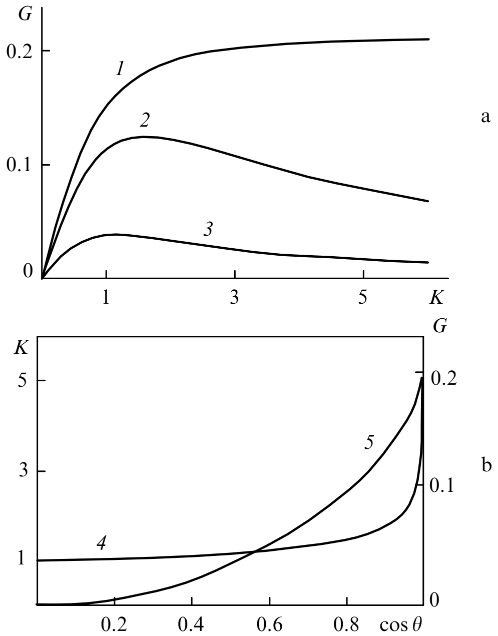

Figure 1. Dependences of the instability growth rate on the wave number k for angles $\theta_0 = 0$ (*1*), π/6 (*2*), π/3 (*3*) in dimensionless variables $G = \gamma/\omega_{\rm pe}$, $K = kc/\omega_{\rm pe}$ (a), and the angular dependences of the optimal wave number k^* (*4*) and the maximum growth rate γ^* (*5*) in the same dimensionless variables (b).

Figure 1b show the dependences of the optimal wave number k^* on $\cos\theta$ [curve (*4*)] and of the maximum growth rate value γ^* on $\cos\theta$ [curve (*5*)]. One can see that the maximum growth rate has a large enough angular width $\Delta\theta \sim 1$, perturbations with the wave numbers $k^* \simeq (1 \div 2)\omega_{\rm pe}/c$ growing most efficiently in a greater part of this angular region. In the angular region $0.1 \lesssim \theta \lesssim 0.8$, the growth rate an be estimated by using the approximate interpolation expression

$$\gamma^* \simeq 0.71\omega_{\rm pe}\frac{v_0}{c}\left(1 - \frac{2.5\theta}{\pi}\right), \tag{7}$$

which approximates with a good accuracy the exact result obtained by solving dispersion equation (5). The maximum growth rate $\gamma^*(\theta)$ is compared with expression (7) for several values of v_0/c in Fig. 2.

If the plasma instability for the direction along $\boldsymbol{n}$ is the transverse Weibel instability [5, 7], then upon deviation from this direction, the potential quasi-static component

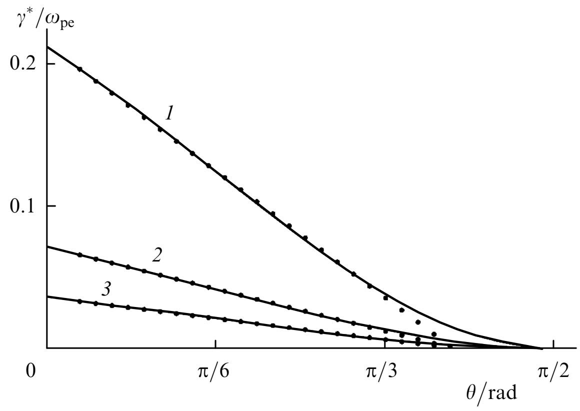

Figure 2. Angular dependence of the maximal (in wave numbers) increment γ^* for $v_0/c = 0.05$ (*1*), 0.1 (*2*), and 0.3 (*3*). The dots show dependence (7).

δn of growing perturbations appears which produces perturbations of the electron density. The relation between the transverse and longitudinal components of the growing low-frequency field characterises the efficiency of excitation of such fluctuations. It is known that the energy of the electromagnetic Weibel instability is mainly contained in the magnetic field; in the geometry under study, this is the component B_y. The longitudinal component of the field is E_z. We will analyse the relation between the longitudinal and transverse components of the electromagnetic field of the instability by using the homogeneous system of linear equations for the electric-field components and the relation between electric and magnetic fields. It is easy to obtain the expression

$$\left|\frac{E_z(\omega,\boldsymbol{k})}{B_y(\omega,\boldsymbol{k})}\right| = \frac{\gamma\,\delta\epsilon_{xz}}{kc(\gamma^2 + \delta\epsilon_{zz})}. \tag{8}$$

In the general case, relation (8) is determined by the wave number and the perturbation propagation direction. However, the characteristic relative value of the fields can be found by estimating the ratio $|E_z^*/B_y^*|$ for the optimal value of the wave number k^*. The dependence of this relation on θ is shown in Fig. 3. In the angular region $\theta \in \{0, \Delta\theta\}$, the longitudinal electric field proves to be of the order of the magnetic field, which demonstrates the efficient excitation of density perturbations during the development of the

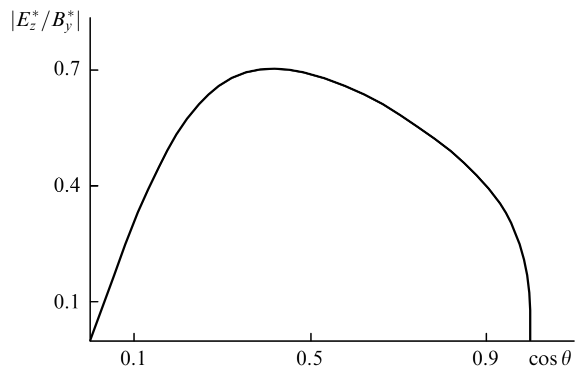

Figure 3. Modulus of the ratio of the longitudinal component E_z^* of the electric field to the magnetic component B_y^*, which corresponds to the optimal wave number k^*, as a function of cosθ.

instability of the anisotropic photoionised plasma. For large angles ($\theta > 70°$), the ratio $|E_z^*/B_y^*|$ considerably decreases; however, the instability for such large angles is suppressed due to the smallness of the growth rate γ.

3. Role of the longitudinal – transverse instability in the interaction of laser radiation with plasma

Our study has shown that a strong longitudinal – transverse instability can develop in a plasma with the anisotropic energy distribution of electrons caused by a short ionising laser pulse. Such instability develops with the growth rate which can amount to a few tenths of the electron plasma frequency. In the case of a low contrast of laser radiation, when a prepulse produces a 'preplasma', this instability develops most efficiently in a plasma corona near a surface reflecting radiation, whereas in the case of high contrast of laser radiation, the instability develops in the skin layer. In the latter case, the approach considered above can be applied for perturbations propagating almost along the target surface (at large angles to the normal), for which the condition of the plasma quasi-homogeneity is fulfilled. Otherwise, the characteristic wave number $k^* = (1 \div 2)\omega_{\rm pe}/c$ is of the order of the inverse thickness ($\omega_{\rm pe}/c$) of the skin layer, and a more complicated method is required to describe the instability. The development of this instability will be accompanied by the exponential growth of fluctuations of the magnetic field and plasma density. In the latter case, $\langle \delta n^2 \rangle_{\boldsymbol k} \sim \langle \delta n^2 \rangle_{\boldsymbol k}^0 \exp(2\gamma t)$, where $\langle \delta n^2 \rangle_{\boldsymbol k}^0$ is the level of initial perturbations of the plasma density corresponding to the spontaneous thermal noise $\langle \delta n^2 \rangle_{\boldsymbol k}^0 \approx n_{\rm e} k^2 r_{\rm D}^2/(1 + k^2 r_{\rm D}^2)$, and $r_{\rm D} = v_0/\omega_{\rm pe}$ is the Debye radius [16].

It is known that non-thermal small-scale fluctuations of the electron density can lead to the additional absorption of laser radiation [12, 17], which is described by the effective collision frequency $\nu_{\rm eff} \sim \omega_{\rm pe}(\delta n/n_{\rm e})^2$. In this connection this instability, which enhances the density fluctuations, can be considered as the reason for increasing the absorption of laser pulses under the condition that the frequency ν of electron collisions producing an isotropic plasma is small compared to the instability growth rate. Note that the artificial microstructuring of target surfaces also leads to the increase in absorption [13 – 15]. The production of a plasma with small-scale density inhomogeneities can enhance the development of the instability because it will be excited from the initial level of density perturbations exceeding the thermal level. Note that the laser plasma is capable of such a self-structuring. Thus, the picture of a standing laser wave (Fig. 4) produced in the focal spot by a laser beam incident at an angle on a target represents a sequence of intensity minima and maxima in the direction along the surface. Due to the pondermotive action of radiation, this leads to the appearance of density inhomogeneities along the target surface with the characteristic wave number k comparable to the wave number k_0 of laser radiation

$$k = 2k_0 \sin\theta_0, \tag{9}$$

where θ_0 is the angle of incidence of the laser beam (for the geometry used in section $\theta_0 = \pi/2 - \theta$). Under the resonance condition $k \simeq k^*$, the initial perturbations of the electron density specified by a standing laser wave will be

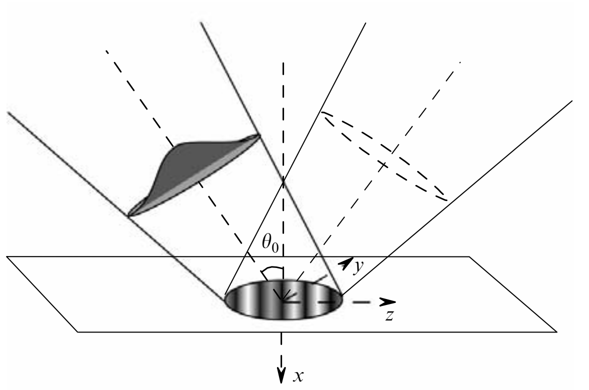

Figure 4. Scheme of the formation of a longitudinal density inhomogeneity along the target surface upon the oblique incidence of a laser beam on the plane target surface.

efficiently enhanced, which can lead to considerable collisionless absorption of laser radiation.

In the case of a high contrast of laser radiation interacting with a material of the solid-state density, the instability can be efficiently enhanced only at the high intensity of laser radiation. Indeed, although the tunnelling ionisation of the surface of a solid target appears at comparatively low laser pulse intensities $I_0 \gtrsim 10^{14}$ W cm^{-2}, the electron energy ($\gtrsim 20$ eV, $v_0/c \gtrsim 0.01$) is insufficient for neglecting the influence of Coulomb collisions of electrons on the instability enhancement according to the requirement $\gamma^* \gg \nu$. However, for $I_0 \gtrsim 10^{17}$ W cm^{-2}, the instability becomes possible even in the case of the solid-state density of the plasma considerably exceeding the critical value. The instability develops with the large growth rate $\gamma^* \gtrsim 0.1\omega_{\rm pe}$ and therefore the additional increase in the instability growth rate due to the pondermotive modulation of the electron density along the target surface [because resonance condition (9) for $k = k^*$ is not satisfied] is not critical.

If the laser prepulse intensity is sufficient for producing a preplasma (corona) near the target surface, the main laser pulse can create the anisotropic distribution of electrons in the corona, by ionising not atoms but ions (as a rule, with small charges) in the preplasma. Taking into account that laser radiation penetrates inside the corona to the cut-off surface with the concentration $n_{\rm e} = n_{\max} \equiv n_{\rm c} \times \sin^2\theta_0$ ($n_{\rm c}$ is the critical plasma density), the instability growth rate in such rarefied plasma for $\theta_0 \sim 1$ is $\gamma^* \sim 0.3\omega_0(v_0/c)$. Because the anisotropy appears only due to the ionisation of ions with small charges, which already exist in the plasma, the laser radiation intensity should be considerably higher than upon photoionisation of neutral atoms, $I_0 \gtrsim 10^{15} - 10^{16}$ W cm^{-2} ($v_0/c > 0.05$). In this case, the electron collision frequency in the cut-off region $n_{\rm e} \simeq n_{\max}$ is negligibly small compared to the characteristic instability growth rate. In addition, according to (9), the modulation of the electron density in this region can satisfy the resonance condition because $k^* \simeq (1 \div 2)\omega_{\rm pe}/c \simeq (1 \div 2)\omega_0 \sin\theta_0/c \approx 2k_0 \sin\theta_0$, which can enhance the instability due to large initial perturbations of the electron density.

4. Conclusions

We have called attention to the important role of a longitudinal – transverse instability which can appear in the anisotropic plasma produced due to the tunnelling ionisation of substance by a short laser pulse. We have described in detail such aperiodic longitudinal – transverse instability. It has been shown that, unlike the electromagnetic Weibel instability discussed in [10, 11, 14], in the case of perturbations propagating at arbitrary angles, electrostatic perturbations increase in the plasma with the anisotropic electron distribution along with magnetic-field perturbations, the potential component of the increasing field being comparable with the transverse component in a broad angular region.

The values of the typical growth rate and wave number describing the properties of such instability have been obtained. It has been shown that the instability is characterised by a broad growth rate over the angle, $\Delta\theta_0 \sim 1$. Unlike the Weibel instability, the longitudinal – transverse instability in the strongly anisotropic plasma develops in a narrower region of wave numbers. The 'optimal' wave number in a broad angular region of the instability enhancement weakly depends on the propagation angle of perturbations and is $k^* \simeq (1 \div 2)\omega_{\rm pe}/c$. The value of the corresponding growth rate is determined by expression (7) and can be of the order of one tenth of the electron plasma frequency.

Small-scale fluctuations of the electron density appearing due to the instability development can cause strong absorption of laser radiation. Simultaneously, the scattering of laser radiation by these electron-density fluctuations can be increased. This should be taken into account in the interpretation of experiments and numerical simulations of the interaction of short laser pulses with plasma. The latter has been based so far on the method of 'large particles' neglecting the tunnelling ionisation of substance by the laser pulse. Such a consideration under certain condition can be inadequate to experiments.

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