PARTICLE ACCELERATION IN PLASMA

FORMATION OF LASER PREPLASMAS TO CONTROL PARTICLE ACCELERATION EFFICIENCY

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Abstract. The issues of modeling the target expansion under the action of a nanosecond laser pulse are considered in order to characterize the plasma torch on the irradiated side and study the possibility of its use for efficient acceleration of charged particles by a powerful short laser pulse. It is shown how various physical models embedded in hydrodynamic calculations affect the modeling results.

Keywords: *particle acceleration, preplasma, hydrodynamic modeling, equations of state*

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1. INTRODUCTION

Short laser pulses with relativistic intensity allow to effectively accelerate electrons and ions of a plasma target to high energies, which opens up prospects for the creation of compact charged particle accelerators with a wide range of practical applications [1-4]. The energies gained and the number of accelerated particles are determined, in addition to the parameters of the laser pulse itself, by the characteristics of the targets used, their structure and density. To obtain electron beams with maximum energies, it is most advantageous to use low-density gases, which allow to accelerate relatively small-charge (typically at the level of tens to hundreds of picocoulombs) electron beams to energies of up to 10 GeV [5]. An increase in the number of electrons with high energies is possible by using denser targets with densities of the order of the critical density (for a wavelength of 1 μm, the critical density is $1.1 \times 10^{21} \text{cm}^{-3}$). It is precisely these targets that make it possible to achieve record values of laser radiation energy conversion into accelerated electron energy [6, 7]. At the same time, the creation of targets with the required optimal density and dimensions for the most effective acceleration of electrons (and then ions) still appears to be a complex task, for the solution of which aerogels, pre-homogenized foam targets [6], cluster and structured targets are used.

One of the simplest ways to manipulate the properties of a target is to create an extended preplasma on the irradiated side of the target with a nanosecond prepulse that precedes the main short pulse, or a synchronized additional pulse. Despite the high contrast of modern laser systems (it reaches values$\sim 10^{10}$) the peak intensity is so high that even with such a contrast the flows in the pre-pulse are sufficient to form plasma, which can increase the efficiency of particle acceleration. Sometimes it turns out that reduced contrast without significant changes in peak intensity values or an optimal delay between the main and additional pulses [8] leads to more efficient plasma formation and, accordingly, electron acceleration. There are experimental works predicting the existence of an optimal gradient at the front of the irradiated target for ion acceleration [9]. Thus, the plasma torch formed under the action of a nanosecond pulse/pre-pulse allows for more efficient acceleration of charged particles, and modeling its formation is a necessary component of the optimization of the acceleration process.

As a rule, the target expansion under the action of nanosecond laser pulses is described within the framework of the hydrodynamic approach. In this case, the laser is an external energy source that is entirely responsible for the creation of the plasma corona: it forms and heats the plasma, increases the pressure in it, the gradient of which causes the target expansion. Several physical effects play a major role at this stage: collisional heat transfer, ionization, and the properties of matter. Within the framework of the hydrodynamic approach, the latter is described using the equation of state and the radiation properties of matter (radiation ranges). At relatively low intensities ($\lesssim 10^{13}$W/cm^2), which are considered in this paper, radiation transfer does not play a significant role, so it can be neglected. The equation of state, on the other hand, is decisive in describing the dynamics of target expansion. A wide-range equation of state is required, since the substance passes through various states during laser irradiation. An initially cold target at room temperature has a normal solid density. This state is preserved for a part of the target even after the start of laser irradiation, which evaporates the irradiated part of the target, forming a hot region of low density, a plasma laser corona heated by a laser pulse or a heat flux from the laser absorption region. This hot region of low density is well described by the equation of state of an ideal plasma (taking into account the variable degree of ionization). Accordingly, there is also a transition region, which has an average density and is also heated. The paper will consider two versions of the equation of state: a wide-range one and one for an ideal plasma. Comparison of the calculation results with them will show what inaccuracies result from using the equation of state, applicable primarily to the plasma corona. The formation of preplasma from a nanosecond pulse, which was then replaced by pico and femtosecond pulses, was considered in [10, 11] in a one-dimensional approximation. In this paper, the simulation is carried out in RZ geometry with symmetry relative to the axis of propagation of the laser pulse, which allows for full consideration of multidimensional hydrodynamic effects.

In addition to the properties of substances, the formation of preplasma is also affected by the efficiency of absorption of the laser pulse, which depends on its intensity and duration. The interaction of laser radiation with a substance is determined by the permittivity, in which the contributions of individual effects depend on the state of the substance. In this paper, a nanosecond laser pulse with an intensity of about or more is considered 10^{12} W/cm^2. At such high intensities, already at the pulse front on picosecond scales, rapid plasma formation by a propagating thermal wave occurs, and the main laser radiation interacts with a hot plasma target, the absorption in which is determined by inverse bremsstrahlung heating during collisions of electrons with ions. Of course, at the initial stage of irradiation, when the structure of the substance is preserved, it is important to take into account both electron-phonon collisions [12, 13] and an accurate description of the transitions between different phases of the substance, including possible metastable states [14]. Moreover, a correct description of the permittivity in a wide temperature range for different states of the target substance is decisive for modeling the effect of femtosecond and even picosecond laser pulses on the target [15, 16]. For such short pulses, additional effects associated with the nonequilibrium distribution of electrons due to ionization [17-20] can arise, leading to features of the absorption of laser radiation. However, for the nanosecond laser pulse considered in the work, the main formation of the plasma corona occurs at times exceeding the characteristic collision times, when the electrons and ions of the expanding plasma are in equilibrium states, and the processes at short times overlap in influence with plasma processes. In this case, a rough estimate of the effect of absorption at the initial stage of interaction was made by a small modification of the collision frequency in the absorption model at low temperatures.

One of the goals of this paper is to compare the dynamics of target expansion for different equations of state, which allows us to show the importance of taking into account the cold component of the equation of state. The paper verifies how different approximations (one-dimensionality of calculations, simplified equation of state) affect the properties of the forming preplasma. The paper also studies the effect of a short burst of laser intensity, occurring at the stage of existence of a developed plasma corona, on its further dynamics. This simple formulation models the effect of a rapid increase in intensity in femtosecond pulses.

2. MODEL FOR CALCULATING PREPLASMA FORMATION

The formation of preplasma occurs at characteristic times of thermal expansion of plasma – on the order of several nanoseconds. Such times significantly exceed the time of collisions between particles, so the dynamics can be described by a hydrodynamic model. An exception is the energy exchange between electrons and ions in a low-density medium that occurs at times on the order of the collision time, so the model must take into account the difference in their temperatures. This effect is enhanced by the fact that laser radiation heats up the electron component. Part of the laser radiation is reflected, the rest is absorbed and subsequently redistributed using a heat flow, which requires taking into account heat transfer (the dominant mechanism is electron heat transfer). The system of equations has the form

$$\partial_t \rho + \partial_j(\rho v_j) = 0, \qquad (1)$$

$$\partial_t(\rho v_k) + \partial_j(\rho v_k v_j) + \partial_k(p_e + p_i) = 0, \qquad (2)$$

$$\partial_t(\rho e_e) + \partial_j(\rho v_j e_e) + p_e \partial_j v_j = Q_{as} + Q_{ei} - \partial_j q_j^{(e)}, \qquad (3)$$

$$\partial_t(\rho e_i) + \partial_j(\rho v_j e_i) + p_i \partial_j v_j = -Q_{ei}. \qquad (4)$$

Here ρ, v_i — density and velocity of the medium, $p_{e,i}, e_{e,i}, T_{e,i}$ — pressure, internal energy and temperature of electrons (ions), respectively; Q_{as} - laser energy release, Q_{ei} — collisional exchange of energy between electrons and ions, $q^{(e)}$ — electron heat flux. Details of the hydrodynamic model can be found in [21]. The calculations use a modified classical thermal conductivity model with a thermal conductivity coefficient corresponding to a hot plasma [22], with the introduction of a limitation on the heat flux corresponding to the free movement of electrons with a limitation coefficient equal to $f = 0.15$. It was verified that for the parameters under consideration the limitation coefficient does not affect the obtained results. At the same time, the presence of thermal conductivity is important for energy transfer from the laser absorption region to the dense region of the target. The electron-ion exchange coefficient is determined by electron-ion collisions [22]. In the laser absorption region, throughout the entire action of laser radiation (several nanoseconds), a difference in electron and ion temperatures is observed, so taking this effect into account is necessary in our model. In the region with densities above the critical density, the collision frequency increases (due to the increase in density and the drop in temperature) and the temperatures of electrons and ions equalize.

This system is solved in the author's multidimensional numerical code FRONT, developed for plasma physics problems. Calculations are performed on the Euler grid, the numerical scheme for the hydrodynamic equations is based on the Godunov-type scheme. To take into account additional physical effects, splitting by physical processes is used, exchange terms and heat transfer are calculated using a completely implicit numerical scheme, which allows for stable calculations at any electron-ion collision times and thermal conductivity coefficients. The code is well parallelized, but the calculations presented below are not particularly resource-intensive and require about 200 processor hours.

In the calculations below, we will assume that the laser radiation falls on the target at a normal angle, and we will also neglect the effects of refraction. In this case, the problem of radiation propagation is simplified - the rays move along a straight trajectory to the critical electron density and then in the opposite direction. The main energy release occurs near the critical density, so this approximation is acceptable. To describe the absorption and propagation of laser radiation, an equation for the intensity is solved, which is integrated along the beam trajectory

$$\frac{dI}{dl} = -kI\,. \quad (5)$$

Here I — intensity of laser radiation in the beam, k — absorption coefficient. Primary ionization of the medium and, accordingly, absorption are determined by multiphoton processes and are described by the Keldysh theory [23]. However, ionization by the field occurs at times less than several picoseconds, which constitute a small part of the pulse duration, and the losses of laser radiation due to ionization are negligibly small, which allows this effect to be neglected. The main absorption mechanism at the intensities under consideration 10^{12}-10^{13} W/cm^2 is the reverse braking absorption, which determines the type of coefficient k [12, 13, 24], calculated using the imaginary part of the permittivity, for which the Drude model is used.

Let's consider the features of the target's equations of state. Due to the two-temperature model, electrons and ions require their own equation of state. First, let's consider the simplest version of the equation of state - an ideal plasma with a variable charge composition. In this case, the pressure is proportional to the concentration of particles of each type

$$p_e = n_e k_B T_e, \quad p_i = n_i k_B T_i. \quad (6)$$

According to equations (2)–(4), the motion of the medium is determined by the gradient of the total pressure, and individual pressure components are needed to calculate the change in the internal energy of the corresponding component. The energy density in the case of an ideal plasma has the form: $\rho e_e = 1.5 n_e k_B T_e, \rho e_i = 1.5 n_i k_B T_i$. The main ionization processes under the conditions under consideration are associated with collisions, therefore the ionization equilibrium is described by the Saha equation [25]. The total ion concentration has the form

$$n_i = \sum_{\alpha} \sum_{\beta=0}^{Z_{\alpha,max}} n_{\alpha,\beta}, \quad (7)$$

Where $n_{\alpha,\beta}$ — concentration of ions with charge β element type α ($Z_{\alpha,max}$ — the charge of the element's nucleus), and the sum is taken over all elements and ion charges. The electron concentration is written as

$$n_e = \sum_{\alpha} \sum_{\beta=1}^{Z_{\alpha,max}} \beta n_{\alpha,\beta}. \quad (8)$$

As noted above, such an equation of state works well for low-density plasma, i.e. with a density significantly lower than that of a solid. ρ_0 for a given substance (for aluminum $\rho_0 = 2.7$g/cm^3). At the same time, it does not allow to describe the behavioral features of a substance near ρ_0, which leads to erroneous dynamics of the solid target.

To construct an equation of state applicable to both the solid phase and the low-density corona, we use the approach proposed in [26]. The electron component is described using the Thomas–Fermi model, which takes into account the degeneracy of electrons at high density and also passes into the ideal plasma model at low density and high temperature. The ionization state is calculated in this model in a consistent manner with the solution of the electron distribution problem. Its value is close to the value given by the Saha model. For the ion component, the Cowan model is used, which offers interpolation expressions for the free energy of ions taking into account three phases: solid, liquid, and gas. The transition between these phases is specified using the melting point and the Debye temperature, for which empirical dependences on the density are also presented. The ion pressure does not depend on the ionization state, so this interpolation is sufficient. These two models are supplemented by a semi-empirical correction for the binding energy, which allows one to obtain the atmospheric pressure at normal substance density and room temperature. The correction is determined by the bulk compression modulus under normal conditions. This model describes the shock-wave compression of substances well, which is confirmed by a comparison with experimental data. In addition, at low densities and high temperatures, it predicts results close to the model of an ideal plasma with variable ionization (see Fig. 1), so it should also correctly describe the expansion of plasma at high temperatures (which occurs with laser irradiation). This equation of state is numerically implemented as a module to the hydrodynamic code. With its help, tabular data are calculated, which are then used in calculations. The following parameters are used to construct the equation of state of aluminum: $Z = 13, A = 26.98$, normal density of matter $\rho_0 = 2.7$ g/cm^3, bulk compression modulus $B = 76$ GPa. Such an equation of state taking into account the degeneracy effects will be the main one in the calculations (let us call it, as the authors of the article [26], QEOS). The total pressure isotherms for such a model in comparison with the ideal plasma model are presented in Fig. 1.

3. CALCULATIONS OF TARGET DYNAMICS UNDER THE ACTION OF A NANOSECOND LASER PULSE

Let's consider the spread of an aluminum plate with a thickness $h = 6$ μm, which is irradiated by a laser pulse with a constant intensity over time (Fig. 2). The characteristic time of intensity increase in the calculations is 20 ps, and it reaches a plateau at the time of 100 ps. The pulse duration in all calculations is $\tau = 3$ ns. The laser radiation falls along the normal to the target surface. Due to the symmetry of the problem, the calculation is carried out in cylindrical geometry RZ (with angular symmetry ϕ), which allows us to fully take into account the three-dimensional expansion of plasma.

Let's define it as Z axis along which the laser radiation propagates. Since we are considering the dynamics created by the prepulse of a tightly focused short pulse, the prepulse focusing radius remains small. Thus, we define the spatial profile of the incoming laser pulse as

$$I(r) = I_0 \exp\left(-\frac{r^2}{r_l^2}\right), \quad r^2 = x^2 + y^2. \tag{9}$$

Here I_0 — intensity at the center of the pulse, r_l — focusing radius, in calculations we will use the value $r_l = 4$ μm. The numerical code used does not allow working with vacuum states, so the area around the target is filled with a low-density substance $\rho_{low} = 10^{-5}$ g/cm^3 (with the same equation of state as is used to describe the target). This value is small enough that its value does not affect the further dynamics of the system. At the initial moment of time, a normal uniform pressure (1 atm) is set in the entire computational domain, which corresponds to the normal density of aluminum and the initial room temperature for the used equation of state QEOS.

According to the Thomas-Fermi model, the average charge of aluminum under normal conditions is $Z = 2.4$. With such an average charge, the electron concentration is formally higher than the critical value, which leads to absorption of laser radiation through the reverse braking mechanism on the front surface of the target. The substance heated at the boundary quickly turns into plasma and flies away from the target, forming a plasma corona.

Over times of about 100 ps, when the maximum intensity of the irradiated laser pulse reaches the target, the temperature in the absorption region rises to $\sim$100 eV (for nanosecond pulse intensity 10^{12} W/cm^2 the heating time to 100 eV is 150 ps, and for the intensity 10^{13} W/cm^2 — 90 ps). It is this temperature that determines the characteristic speed of expansion of the corona of the formed plasma, $2c_s / (\gamma - 1) \sim 100$ km/s, where c_s — the speed of sound in the corona, proportional to the root of the temperature $c_s \propto T^{1/2}$. Note that after the incident radiation intensity reaches a plateau, the temperature changes slightly. Over times of about 300 ps, the hot plasma has time to fly away towards the laser pulse, filling almost the entire computational domain. Note that the hydrodynamic velocity of expansion obtained in the calculations turns out to be significantly less than the velocity of expansion of collisionless plasma into a vacuum [27], according to which the velocity of the front of the expanding torch increases with time, $v_f \simeq 2c_s \ln(\tau + \sqrt{\tau^2 + 1})$, Where $\tau = \omega_{ph} t / \sqrt{2e}$ and for the simulation parameters $v_f \sim 1500$ km/s.

In the opposite direction, deep into the target, a thermal wave goes, which forms a shock wave. At the intensities under consideration, the first shock wave that runs along the target is relatively weak (at $I = 10^{12}$ W/cm^2 wave speed $D = 7.5$km/s, which is consistent with the Hugoniot relations for ablation pressure ~ 330 kbar). This wave, traveling at a speed comparable to the speed of sound in cold metal, accelerates the plate material. As a result, the plate begins to shift in the region of laser absorption (Fig. 3, 4). A series of weak shock waves that propagate along the plate further are clearly visible in the density distributions (Fig. 5).

The scattered hot plasma fills the space around the target from the front side (from the side of the incident laser pulse). The maximum temperature of several hundred eV is observed near the laser absorption region, the energy removal together with the substance or due to heat transfer maintains the temperature in the plasma corona. The heat flow also exists in the direction of the dense area of the target, but due to the high heat capacity, the dense parts of the target remain relatively cold.

Subsequently, the accelerated substance of the plate in the region near the laser absorption will continue its movement, the target gradually becomes thinner and eventually the plate breaks through (Fig. 4). The characteristic time of complete burnout of the target with a laser pulse duration of $\tau = 3$ ns and for intensity $I = 10^{12}$ W/cm^2 makes up $t_p = 15$ns, for intensity $I = 10^{13}$ W/cm^2— $t_p = 6$ ns, and for intensity $I = 5 \times 10^{13}$ W/cm^2— $t_p = 3.5$ ns.

It should be noted that the described dynamics of the plate is possible only due to the presence of a cold component in the equation of state, which allows maintaining the elasticity of the dense medium. This is clearly seen when comparing the calculations performed with calculations with another equation of state - an ideal plasma, which does not contain degeneracy. In the equation of state of an ideal plasma, due to the absence of a cold part, the speed of sound in the medium is determined by the temperature $c_s \propto T^{1/2}$. In a dense medium at the initial moment of time the initial low temperature of the target $T_0 = 10^{-2}$ eV gives the speed of sound $\sim$200 m/s, which is significantly lower than the speed of sound in cold aluminum at normal density. The speed of sound characterizes the speed of propagation of small disturbances in such a medium. As a result, in the calculations the plate does not shift entirely, as shown above, but is raked by the first shock wave propagating from the point of energy release (Fig. 6).

At the same time, it is worth noting that the parameters of the low-density region - the laser corona for both equations of state are very close to each other. Figure 7 shows a comparison of the electron density and temperature profiles along the straight line $r = 0$ (at the center of the laser beam) for two equations of state at a time of 2 ns. It is seen that the profiles n_e And T_e in the low-density region are close for both calculations. We define the density gradient as $\Gamma = n_e^{-1} dn_e / dz$. For calculation with the QEOS equation of state $\Gamma_{QEOS} = 2.5$ microns^{-1}, and for calculation with the equation of state of an ideal plasma $\Gamma_{IP} = 3.1$ microns^{-1}, which gives a difference in $\sim 20\%$. The obtained agreement is explained by the fact that the profiles in the low-density region are determined only by the properties of the substance in such states and the magnitude of the energy release. A more noticeable difference arises for the position of the critical density, which is spatially shifted by several microns into the target when using the equation of state of an ideal gas. The position of the critical region is important for determining the focal point of a short laser pulse, since the efficiency of particle acceleration depends on it. In general, if we are interested only in the properties of the laser corona, the equation of state of an ideal plasma is sufficient.

For a more detailed description of the dynamics, let us consider the pressure behavior (Fig. 7). In the low-density corona region, it coincides for both calculations (up to an offset, which is consistent with the density behavior). In the high-density plasma region, the difference in pressure behavior explains the difference in the target dynamics. For the equation of state of an ideal plasma, the peak pressure value is located at the boundary with the dense plasma, and the pressure drops inside the layer. This configuration leads to the effect of matter raking. For the second equation of state (with degeneracy), the maximum pressure value is in the high-density region and is associated with the shock wave that passed through the medium, compressed, and accelerated the target. It is worth noting two points at which the pressure goes negative. These points are in the region of transition from a dense to a low-density medium at low temperatures: in this case, the state of the medium is in the region of non-monotonic pressure behavior according to the QEOS predictions (see Fig. 1). This region requires a more physically complex model for constructing the equation of state. The small volume in space occupied by regions with such a state of matter does not affect the large-scale dynamics of the target, but the presence of these regions indicates the need for further refinement of the equation of state.

Fig. 8 shows a comparison of the target shape as a result of irradiation with pulses of different intensities and the same focusing radius.r_0. Due to the difference in temperature in the corona and, accordingly, pressure, the displacement of the plate is different, as is the radius of the hole formed in the plate. It is clearly seen that at high intensities the target bends, and as a result the dense regions that surround the expanding plasma begin to influence the flow. This effect cannot be taken into account within the one-dimensional approximation, which demonstrates the importance of taking into account the three-dimensional expansion, especially under conditions of a small focusing radius of the laser pulse.

Note that simpler one-dimensional calculations are often used to estimate the size and characteristic temperature gradients of an expanding plasma cloud, sometimes using approximate renormalizations for the expansion of the plasma cloud in transverse directions, for example, using an additional dynamic equation [28].

Fig. 9 shows a comparison of the electron density and temperature profiles for three-dimensional (RZ) and one-dimensional (1D) calculations. The difference is especially noticeable in the electron density profiles. In the one-dimensional calculation, the corona is strongly elongated with a gentle gradient (for $I_0 = 10^{12}$W/cm^{2}gradient $\Gamma_{1D} = 0.2$ microns^{-1}, which is an order of magnitude lower than inRZcalculation). In three dimensions (RZ) the calculation clearly shows the presence of several characteristic gradients - at densities several units higher than the critical density, near the critical density and in the low-density plasma corona. The plate displacements also differ significantly: in the one-dimensional case, the plate will shift several times more strongly.

The inverse bremsstrahlung absorption model used in the calculations does not quite correctly describe the absorption behavior at low temperatures. The absorption coefficient is determined through the imaginary part of the plasma permittivity, $k = 2\omega_l \operatorname{Im}\sqrt{\varepsilon} / c$ (Here ω_l— the frequency of the laser radiation) and is proportional to the effective collision frequency. For the plasma model of electron-ion collisions used in the code, the collision frequency increases with decreasing temperature, $\nu_{ei} \propto T_e^{-3/2}$, which also leads to an incorrect increase in absorption at low temperatures [29]. As a simple test of how much such enhanced absorption plays a role, we will use a simple modification of the presented model: when calculating the electron-ion frequency, we will limit the temperature from below by the Fermi temperature for the substance:$\nu_{ei} = \nu_{ei}(\max(T_F, T_e))$. For aluminum of normal density we obtain (taking into account the equation of state used, which yields the value n_e) $T_F = \hbar^2(3\pi n_e)^{2/3} / (2m_e k_B) = 4.7$ eV. Calculation with such a collision frequency limitation leads to a delay in the growth of the corona temperature, but at very short times $\lesssim 50$ps when is the plasma temperature $T \sim 1$eV. As the temperature increases to tens of eV (at times ~ 100ps) the difference between the calculations disappears. Thus, the correct absorption coefficient at the initial stage of interaction has an insignificant effect on the subsequent dynamics of the target at the energy flows under consideration. This is partly due to the fact that the plasma formation processes occur at low laser intensities compared to the peak ones and affect only a small part of the pulse energy, while the dynamics of the target itself and the laser corona are determined by subsequent stages of interaction, at which a significant part of the laser pulse energy enters the target.

Let us consider the effect of a short picosecond intensity burst on the target dynamics against the background of a constant nanosecond pulse. Such bursts are observed in the pre-pulse of some laser systems [30], and also describe the increase in intensity on the picosecond scale before the arrival of a femtosecond pulse [31]. We will assume that the picosecond pulse has a Gaussian time profile with a characteristic scale $t_1 = 2$ps and intensity $I_1 = 10^{15}$W/cm^2, and the intensity of the nanosecond pulse, duration $\tau_0 = 3$ns, is $I_0 = 10^{12}$W/cm^2. The center of the short picosecond pulse is located 2 ns after the start of the long pulse (see Fig. 8a). Thus, the energy of the short pulse in $\sim I_1\tau_1 / I_0\tau_0 \approx 7$ times greater than the energy of a nanosecond pulse, and it has a significant impact on the dynamics of the target.

The results of calculations with such a pulse are presented in Fig. 10 b and Fig. 11. The results show that at the moment of arrival of the maximum of the picosecond pulse, $t = 2$ns, a short pulse manages to significantly heat the plasma and change the electron density gradient near the critical point: the profile steepens and the gradient Γ changes in a short time from 2.5 μm^{-1} up to 3.2 μm^{-1}. The target itself (its denser part) does not have time to respond to the short pulse at picosecond times. The gradient of the low-density part of the plasma torch also remains virtually unchanged. Thus, the picosecond prepulse has little effect on the target density characteristics, only changing the characteristic gradient near the critical density. At the same time, the presence of such a relatively powerful intensity burst in the prepulse can significantly change the parameters of the expanding plasma at later times, leading to a density burst in the plasma corona (see Fig. 9c). In addition, the rapid energy deposition and the accompanying plasma heating lead to an impact effect, as a result of which a strong shock wave subsequently propagates along the target, which leads to the destruction of the rear boundary of the plate, see Fig. 8b. It should be noted that the observed destruction occurs due to a strong rarefaction wave that occurs after the shock wave is reflected from the boundary. This is evident from the pressure profiles shown in Fig. 9 — a large area with negative pressure appears. Since the calculations did not use a model of destruction or elastic plasticity, the observed effect is associated with the behavior of the equation of state. Therefore, this result is rather an indication of possible destruction, and a more detailed study of this process requires calculations with more accurate models of the equation of state, as well as taking into account the strength properties of substances.

4. DISCUSSION OF RESULTS AND CONCLUSION

The paper examines the process of laser corona formation when a metal (aluminum) target is irradiated with nanosecond pulses and intensities $10^{12} - 5\times10^{13}$ W/cm^2. The interest in such calculations is associated with the need to correctly obtain the parameters of the plasma formed before the arrival of a powerful femtosecond pulse, as well as the ability to control the parameters of this plasma. A physical model required for such calculations is presented, and the main effects accompanying the formation of plasma at different laser intensities are discussed. In particular, the effect of two equations of state on the calculations is considered. The equation of state of an ideal plasma with variable ionization is suitable for describing the low-density region of the plasma crown, in which calculations using it are consistent with the results of calculations using a wide-range equation of state. Thus, if we are interested only in the properties of the low-density corona, the use of such an equation of state is justified. To correctly describe the dynamics of the dense part of the target, as well as its displacement, equations of state are required that take into account the effect of electron degeneracy and bonds between atoms at high matter density. It is also shown that within the framework of a one-dimensional model, erroneous gradients of the electron density in the region of the critical density are obtained: due to the effects of lateral expansion of the plasma, the gradients become steeper. This is especially important in the case of a small focusing radius, which exactly corresponds to the pre-pulse of short laser systems. The presented model allows calculating the dynamics of thin plates, their deflection and displacement, which is important to consider when focusing laser radiation.

The target preplasma density profiles predicted by the hydrodynamic calculations are presented in two graphs (Fig. 12): along the axis $Z(r = 0)$ and along the axis R ($z = 11$ μm, i.e. along the beam, which is 5 μm away from the original surface of the plate). From them it is evident that the characteristic gradient in the region of critical electron density is $\Gamma = 2.5$ microns^{-1} and has a weak dependence on the intensity and time of irradiation. Large the difference is observed at densities several units higher than the critical one ($1–10n_c$): the higher the intensity, the more developed the corona is and the flatter the density profile is formed in this region. The characteristic gradient in the region of the low-density corona (at densities from $0.05n_c$ to $0.5n_c$) varies from 0.03 to 0.13 μm^{-1}. Thus, the formed longitudinal profile of the preplasma density has a rather sharp gradient in the region of critical density and an extended preplasma starting from the characteristic density of $0.3–0.4n_c$, the value of which decreases by an order of magnitude over 10–15 μm. This indicates rather limited possibilities for controlling the parameters of the preplasma using a prepulse of spontaneous amplification of the emission of the main powerful short pulse, since in this case, with a fixed delay, only the longitudinal profile of the preplasma can be used, changing slightly depending on the intensity of the prepulse.

Significantly more possibilities for controlling the parameters of the plasma target arise when using an independent nanosecond pulse, especially with the possibility of focusing the main pulse used to accelerate electrons at different angles to the target, and, in particular, almost parallel to the expanding target. In this case, the pulse can propagate along the transverse profile of the formed preplasma ($z =$const), the electron density of which increases with the intensity of the nanosecond pulse. At the same time, close to the target axis, the concentration is practically constant at the level of $0.3–0.4n_c$, when moving away along the radius, it first begins to grow, reaching several critical densities, and then the growth is replaced by a fall. The non-monotonicity of these density profiles (maximum in the region of 15–20 μm from the focusing axis) is associated with the transverse features of the target expansion: outside the radius of laser irradiation, the substance is raked into a protrusion that appears on the front side (Fig. 6). Thus, at sufficiently high intensities of the nanosecond pulse ($\gtrsim 5 \times 10^{13}$ W/cm^2) in front of the irradiated part of the target at some distance from the focal center, a region of near-critical density appears, 5–10 μm in length, which can be used for more efficient electron acceleration. It also seems possible to use the transverse profile of the plasma behind the burnt-out target for efficient electron acceleration. Fig. 13 shows the transverse profiles of the concentration n_e behind the plate at the time near and after the plate breakthrough. It is evident that the electron density decreases with time and is a fraction of the critical density. At the same time, at earlier times, the density in the central part will be higher (of the order of critical), but will be limited by the "walls" of the flying apart target. Thus, by changing the delay between the additional nanosecond and main femtosecond laser pulses, it is possible to achieve interaction of the latter with the most optimal profile of the plasma target. For example, when irradiating an aluminum target 6 μm thick with a nanosecond pulse (3 ns long) with an intensity of 10^{14}W/cm^2, focused into a spot 4 μm in size, 4 ns after its arrival at the target, a homogeneous plasma with a density of the order of $0.2n_c$ and a size of about 100 μm, which is optimal for accelerating electrons in the self-capture mode of a relativistic laser pulse with a duration of 10 fs and an energy of about 2 J [7]. Considering that the characteristic times of a significant change in the density profile of the expanding plasma are hundreds of picoseconds, it is possible to ensure the necessary synchronization of a nanosecond pulse, creating a plasma with a given profile, with a femtosecond laser pulse, for the most effective acceleration of electrons.

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Captions to the figures

Fig. 1. Comparison of total pressure isotherms for the QEOS equation of state (solid curves) and ideal plasma (dashed curves) for aluminum.

Fig. 2. Calculation geometry and time profile of the laser pulse.

Fig. 3. Density distribution (in g/cm^3) and temperatures (in eV) in the target and surrounding space. Calculation with $I_0 = 10^{13}$ W/cm^2. Moment in time$t = 2$ns.

Fig. 4. Plate dynamics under constant intensity pulse irradiation 10^{13} W/cm^2.

Fig. 5. An enlarged portion of the plate in Fig. 2, where the waves running along it are clearly visible ($I = 10^{13}$W/cm^2).

Fig. 6. Density distribution (in g/cm^3) and temperature (in eV) in the target when calculated with the equation of state of an ideal plasma at $I_0 = 10^{13}$ W/cm^2. Moment in time $t = 2$ns.

Fig. 7. Comparison of electron density (solid curves) and electron temperature (dashed curves) profiles for calculations with different equations of state (a). Plasma pressure (solid) and density (dashed) (b). Intensity $I_0 = 10^{13}$ W/cm^2, moment in time $t = 2$ ns.

Fig. 8. Comparison of target shape for different intensities at the end of the laser pulse$t = 3$ns. The intensity values are indicated under the figures. The calculation is performed using the equation of state with degeneracy (QEOS).

Fig. 9. Comparison of electron density (solid curves) and temperature (dashed curves) profiles for calculations performed in geometryRZ and one-dimensional calculations for two intensity variants at the moment $t = 2$ns.

Fig. 10. Time profile of a laser pulse with a picosecond burst of magnitude $I_1 = 10^{15}$ W/cm^{2}in the background $I_0 = 10^{12}$ W/cm^{2}and characteristic duration $\tau_1 = 2$ps (a) . Distribution of target density after pulse action at a time of 3 ns (b) .

Fig. 11. Distribution of electron concentration (at criticaln_c, solid curves) and temperatures (in eV, dashed curves) along the straight line $r = 0$. Settlements with $I_0 = 10^{12}$ W/cm^2(ns), and calculation with $I_0 = 10^{12}$ W/cm^2 and peak value $I_1 = 10^{15}$ W/cm^{2}in a femtosecond pulse (ns+ps). Time moments are 2 ns (a) and 3 ns (b). The distributions of electron concentration (solid curves) and pressure (dashed curves) at the moment 3 ns for this calculation are shown (c).

Fig. 12. Electron concentration profiles along a straight line $r = 0$ (a) and beam $z = 11$ μm (b) for different moments of time and different intensities. Time is in ns, intensities in W/cm^2.

Fig. 13. Density distribution at the moment $t = 4$ For $I = 10^{14}$ W/cm^2 (a) , the beam along which the one-dimensional profiles are constructed is shown. One-dimensional profiles of the electron concentration for two variants of intensity and time after the plate breakthrough along $z = -20$ μm (b) .