Absorption and reflection of laser radiation by a dispersing high-temperature plasma

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The two-dimensional self-consistent problem of absorption and reflection of incident radiation by a dispersing high-temperature plasma is considered. The process of interaction of the radiation with matter is described within the framework of the Maxwell equations for the light-wave field and the gas-dynamical equations with allowance for the electronic thermal conduction. A numerical solution to the indicated equations is presented for a plane plasma layer in a wide range of flux densities and for different laser-pulse durations. The main attention is given to the investigation of the efficiency with which the laser-radiation energy is injected into the plasma. The computations are carried out under the assumption of a bremsstrahlung mechanism of absorption and with allowance for a number of collective processes.

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1. FORMULATION OF THE PROBLEM

One of the most important tasks of the physics of the interaction of high-power laser radiation with matter is the investigation of the efficiency of injection of the light energy into the target plasma. As applied to the problem of laser thermonuclear fusion (LTF), this problem is directly connected with the determination of the hydrodynamic efficiency in the process of heating and compression of a thermonuclear target, [1,2] on the magnitude of which the efficiency of the laser thermonuclear system as a whole largely depends. [3] The physical processes determining the useful radiation-energy portion that goes directly into the heating and compression of the thermonuclear fuel are the process of absorption and reflection of the light and the process of transport of the evolved heat by means of the mechanism of electronic thermal conduction under conditions of highly developed hydrodynamics. All the indicated processes develop at the stage of formation and dispersion of the "corona" of laser thermonuclear targets. [2] In the presently known theoretical investigations of the hydrodynamics of the interaction of laser radiation with matter, the process of light-energy dissipation was studied within the framework of the transport equation for the radiation flux, or under the assumption of complete absorption of the radiation in the vicinity of the point with the critical density. [2,4,5] However, the correct description of the process of absorption and reflection of the incident radiation in the laser plasma is impossible without the inclusion, in the consideration of the problem, of the Maxwell equations for the light field. This circumstance is connected with the presence of substantial gradients in the density and other gas-dynamical quantities in a plasma flying apart under the action of laser radiation. The formulation of such a problem can be found in, for example, [6]. Such an approach was used in an application to the laser plasma in[7], but without allowance for thermal conduction and the dynamics of the dispersion of the plasma. In[8] it is suggested that only the geometrical-optics approximation should be considered with the condition of total reflection at the critical point. However, such an approximation yields an incorrect result in the vicinity of the critical point, as a result of the inapplicability of geometrical optics. In[8] the authors consider, for example, the analytic solution for a linear layer, and the result is compared with the geometrical-optics approximation. However, first, the comparison was not done quite correctly, since the collision rate was assumed in the solution of the wave equation to be independent of the space coordinate. Second, and this is the main thing, in[8] the authors considered already-formed plasma profiles, their optical thickness being such that the radiation is almost totally absorbed by the plasma. In this case, the contribution of the vicinity of the critical point to the optical thickness can be small. But in the hydrodynamics calculation the plasma profile is formed in the course of the irradiation of the layer, and it is possible to have density gradients such that a substantial part of the radiation is reflected from the layer. In this case the calculation cannot be carried out, using the method proposed in[8]. In the case of large plasma-density gradients, the decisive role in the radiation-absorption process is played by the skin effect, which is taken into account in the rigorous approach to the solution of the Maxwell equations.

Let a monochromatic electromagnetic wave $P = P_0 e^{i\omega z/c}$, where ω is the angular frequency and c is the velocity of light in a vacuum, be incident on a plane plasma layer from $z = -\infty$ (Fig. 1). Since the characteristic times of variation of the density and temperature of the plasma exceed by far the quantity $2\pi/\omega$, we can use the quasi-stationary Maxwell equations. The process of heating and dispersion of the plasma is described by the following system of equations:

$$\frac{\partial}{\partial t}\left(\frac{1}{\rho}\right) = \frac{\partial u}{\partial m}, \quad \frac{\partial z}{\partial t} = u, \tag{1}$$

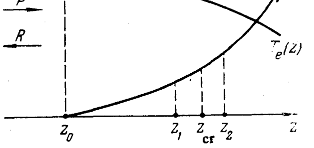

FIG. 1. Density and electron-temperature profiles in the "corona" of the laser target.

TABLE I.

	Radiation-energy flux density q in W/cm²	Electron density, n_e cm⁻³	$\nu_{\rm eff}$ in sec⁻¹
1	$0 \leqslant q \leqslant q_0$	$0 \leqslant n_e \leqslant n_{\rm cr}$	ν_0
2	$q_0 \leqslant q \leqslant \min(q_1, q_2)$	$\left\lvert \left(\frac{n_e}{n_{\rm cr}}\right)^{1/2} - 1 \right\rvert \leqslant \frac{3}{4} \ln^{-1} \frac{M_i}{Zm_e}$	ν_1
3	$q_0 \leqslant q \leqslant \min(q_1, q_2)$	$\left\{1 - \left(\frac{n_e}{n_{\rm cr}}\right)^{1/2}\right\} > \frac{3}{4} \ln^{-1} \frac{M_i}{Zm_e}$	ν_0
4	$\min(q_1, q_2) < q < \infty$	$0.3 \geqslant \left\lvert \left(\frac{n_e}{n_{\rm cr}}\right)^{1/2} - 1 \right\rvert \geqslant \left(\frac{Zm_e}{M_i}\right)^{1/2}$	ν_2
5	$\min(q_1, q_2) < q < \infty$	$\left\lvert \left(\frac{n_e}{n_{\rm cr}}\right)^{1/2} - 1 \right\rvert < \left(\frac{Zm_e}{M_i}\right)^{1/2}$	ν_3
6	$\min(q_1, q_2) < q < \infty$	$\left\{1 - \left(\frac{n_e}{n_{\rm cr}}\right)^{1/2}\right\} > 0.3$	ν_0

$$\frac{\partial u}{\partial t} = -\frac{\partial}{\partial m}\left(p - \mu\rho\frac{\partial u}{\partial m}\right), \tag{2}$$

$$\frac{\partial \mathscr{E}_e}{\partial t} = -p_e\frac{\partial u}{\partial m} + \frac{\partial}{\partial m}\left(\varkappa\rho\frac{\partial T_e}{\partial m}\right) + \frac{1}{\rho}\varepsilon_2\frac{\omega}{8\pi}|E_y|^2 - Q, \tag{3}$$

$$\frac{\partial \mathscr{E}_i}{\partial t} = -\left(p_i - \mu\rho\frac{\partial u}{\partial m}\right)\frac{\partial u}{\partial m} + Q, \tag{4}$$

$$\frac{\partial H_x}{\partial z} + \frac{i\omega}{c}\varepsilon E_y = 0, \tag{5}$$

$$\frac{\partial E_y}{\partial z} + \frac{i\omega}{c}H_x = 0, \tag{6}$$

where u is the velocity, ρ is the density, z is the Euler coordinate, m and t are Lagrangian variables, $p = p_e + p_i$ is the total pressure, $p_e = (\gamma - 1)A_e \rho T_e$ and $p_i = (\gamma - 1)A_i \rho T_i$ are respectively the electronic and ionic pressures, $\gamma = \frac{5}{3}$ is the adiabatic constant, $\mu = \frac{4}{3}\eta_0 T_i^{5/2}$ is the coefficient of ionic viscosity, $\mathscr{E}_e = A_e T_e$ and $\mathscr{E}_i = A_i T_i$ are the internal energies of the electron and ion components, T_e and T_i are the electron and ion temperatures, $\varkappa = \varkappa_0 T_e^{5/2}$ is the coefficient of electronic thermal conductivity, $Q = Q_0 \rho(T_e - T_i)/T_e^{3/2}$ is the rate of exchange of energy between the ions and electrons, E_y and H_x are the complex components of the electric and magnetic fields (below we drop the indices on these fields), and $\varepsilon_1 + i\varepsilon_2$ is the complex permittivity of the plasma. In the case of purely Coulomb scattering of the electrons by the ions, ε is a function of only the density and temperature of the electrons:

$$\varepsilon_1 = 1 - a_0\rho, \quad \varepsilon_2 = b_0\rho^2/T_e^{3/2}, \tag{7}$$

where a_0 and b_0 are known constants. In the case when allowance is made for the anomalous absorption of light by the plasma, [9]

$$\varepsilon_2 = \frac{4\pi e^2 n_e \nu_{\rm eff}(n_e, T_e, q)}{m_e \omega^3}, \quad q = \frac{c}{4\pi}|E|^2, \tag{8}$$

where n_e is the electron density and e and m_e are the electron charge and mass. The effective collision rates, $\nu_{\rm eff}$, are given in Table I, where the threshold fluxes are equal to:

$$q_0 = 5.4\cdot 10^{-18} Z^{3/2}(m_e/M_i)^{1/2} n_e^{3/2}/T_e^{1/2},$$
$$q_1 = 5.2\cdot 10^{15} T_e, \quad q_2 = q_0(\nu_2/\nu_0)^2,$$

and the values of the effective collision rates are given by the following expressions:

$$\nu_0 = 10^{-9} Z n_e / T_e^{3/2}, \quad \nu_1 = \nu_0 (q/q_0)^{1/2},$$
$$\nu_2 = 5.65\cdot 10^4 n_e^{1/2} (Zm_e/M_i)^{1/2}, \quad \nu_3 = 5.65\cdot 10^4 n_e^{1/2} (Zm_e/M_i)^{1/3}$$

(T_e is in keV).

The third term on the right-hand side of Eq. (3) represents Joule heating. For the numerical integration of the system (1)–(6), it is convenient to represent the expression for the energy release in the form of a divergence. It is easy to obtain with the aid of the Poynting theorem the equality

$$\frac{1}{\rho}\varepsilon_2\frac{\omega}{8\pi}|E|^2 = -\frac{\partial \bar{q}}{\partial m}, \tag{9}$$

where $\bar{q} = -(c/16\pi)(EH^* + E^*H)$ is the light-energy flux density averaged over the period of the field oscillations. The system of equations (1)–(6) can be solved for an arbitrary initial distribution of the hydrodynamic quantities over the coordinate z. For the solution of the Maxwell equations, it is sufficient to give the amplitude of the incident wave and use the condition that $E \to 0$ as $z \to +\infty$.

2. SOLUTION OF THE MAXWELL EQUATIONS

The simultaneous solution of Eqs. (1)–(6) amounts to the successive integration of the Maxwell equations for given temperature, $T_e(z)$, and density, $\rho(z)$, distributions. Let us represent the fields E and H in the form

$$E = P + R = P(1 + V), \quad H = \beta(-P + R) = \beta P(-1 + V), \tag{10}$$

where

$$\beta = \sqrt{\varepsilon} = \beta_1 + i\beta_2, \quad \beta_1 = [{}^1/_2(\varepsilon_1^2 + \varepsilon_2^2)^{1/2} + {}^1/_2\varepsilon_1]^{1/2}, \quad \beta_2 = [{}^1/_2(\varepsilon_1^2 + \varepsilon_2^2)^{1/2} - {}^1/_2\varepsilon_1]^{1/2}.$$

The functions P and R coincide respectively with the incident and reflected waves in the vacuum ($z \leqslant z_0$); $V(z \leqslant z_0)$ is the reflection coefficient. The substitution of (10) into (5) and (6) yields

$$\frac{dP}{dz} = \frac{i\omega}{c}\beta P - \frac{1}{2\beta}\frac{d\beta}{dz}[P(1 - V)], \tag{11}$$

$$\frac{dV}{dz} = -2\frac{i\omega}{c}\beta V + \frac{1}{2\beta}\frac{d\beta}{dz}(1 - V^2). \tag{12}$$

Equation (12) in the case of classical absorption contains only one unknown function, V, and can be numerically integrated if the value of V is given at some point. In the case under consideration, the method of computation is connected with the selection of the neighborhood, (z_1, z_2), of the critical point, such that the field E at the point z_2 is negligibly weak, while the geometrical-optics approximation is valid at $z \leqslant z_1$. In practice, the choice of the width of the region (z_1, z_2) is made with the aid of the relation

$$\frac{\omega}{c}\int_{z_1}^{z_{\rm cr}} \sqrt{\varepsilon_1}\,dz = \frac{\omega}{c}\int_{z_{\rm cr}}^{z_2} \sqrt{-\varepsilon_1}\,dz = \alpha\cdot 2\pi, \tag{13}$$

where $\alpha \sim 2 - 3$. In this case $V(z_2) = 0$. Indeed, on account of the rapid attenuation of the field in the region $z > z_{\rm cr}$, we can set $\varepsilon(z) = \varepsilon(z_2)$ for $z \geqslant z_2$. Then the field

E in this region will have the form

$$E=P+R=C_1\exp\left\{i\frac{\omega}{c}\beta_1 z-\frac{\omega}{c}\beta_2 z\right\}+C_2\exp\left\{-i\frac{\omega}{c}\beta_1 z+\frac{\omega}{c}\beta_2 z\right\}. \quad (14)$$

Since $E\to 0$ for $z>z_2$, $C_2=0$.

Integrating Eq. (12) from the point z_2 to $z=z_0$, we find the universal, i.e., incident-field independent, function $V(z)$, after which we can also integrate Eq. (11) for the function P, since $P(z_0)=P_0\exp(i\omega z_0/c)$ where P_0 is the prescribed amplitude of the incident wave. The expression for the flux density $\bar{q}(z)$ for the found functions P and V is of the form

$$\bar{q}=\frac{c}{8\pi}|P|^2[\beta_1(1-|V|^2)+2\beta_2\,\mathrm{Im}\,V]. \quad (15)$$

In the case when allowance is made for the collective effects, which formally amounts to allowance for the dependence of the imaginary part, ε_2, of the permittivity on the square, $|E|^2$, of the field amplitude, the solution is found by the method of iterations. As the initial approximation, we use the solution with a field-independent value of ε_2. The gas-dynamical equations with thermal conduction, (1)–(4), are solved by the well-known difference methods.[10]

3. THE RESULTS

As has already been indicated above, the considered model allows us to give a physically correct self-consistent description of the processes of absorption and reflection of monochromatic radiation by a dispersing plane plasma layer and answer a number of important questions, such as the efficiency of injection of the light energy into the plasma, the spatial structure of the radiation field and, consequently, the composition of the energy release, the shaping of the reflected wave, the temporal evolution of the reflection coefficient as a function of the parameters of the laser pulse and target. Furthermore, the simultaneous solution of the Maxwell equations and the equations of gas dynamics enables us to interpret the experimental data on the reflection with allowance for the actual shape of the laser pulse and for a number of secondary effects, as for example the presence of forepulses. In the present paper the Eqs. (1)–(6) have been solved for neodymium laser radiation ($\omega=1.8\times10^{15}$ sec^{-1}) in the $10^{12}-10^{16}$ – W/cm^2 flux density range and plane, homogeneous (at $t=0$) targets of the type DT, $(CH_2)_n$.

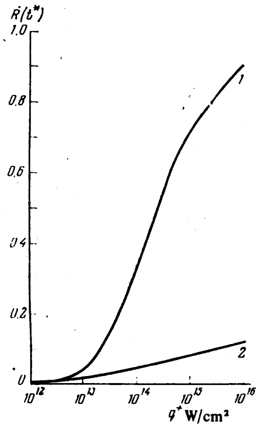

FIG. 2. Dependence of the integrated—over the time interval t^*—coefficient of reflection with respect to energy on the magnitude of the incident flux in the case of classical absorption (the curve 1) and with allowance for anomalous absorption (the curve 2).

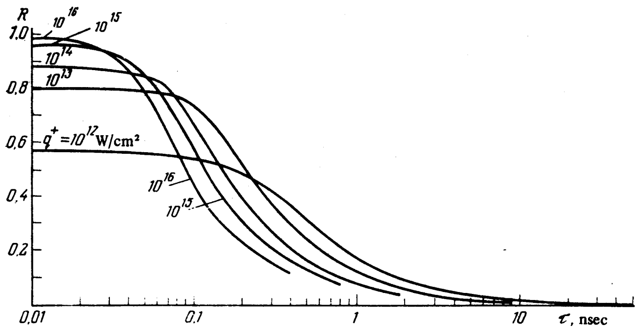

FIG. 3. Dependence of the integrated reflection coefficient on the laser-pulse duration for different values of the incident flux.

In Fig. 2, we show the dependence of the integrated reflection coefficient

$$R=\int_0^{t}q^-\,dt\Big/\int_0^{t}q^+\,dt$$

on the incident flux q^+ for a DT layer ($\rho(0, z)=\rho_0=0.2$ g/cm^2, $T(0, z)=0$). The curve 1 corresponds to classical absorption; the curve 2 was obtained with allowance for anomalous absorption. As follows from the plots shown, at radiation-flux densities $q\gtrsim10^{13}$ W/cm^2 the anomalous-absorption processes begin to play a decisive role, leading to significantly lower reflected-energy values. In the present case the dominant contribution to the absorption in the entire range of incident fluxes q^+ was then connected with the effective frequency ν_1, which corresponds to the process of decay of the light wave into an electronic wave and ionic sound.

In Fig. 3 we show a family of plots of $R(q^+, \tau)$ illustrating the efficiency of absorption, including anomalous absorption, of radiation energy in a DT target as a function of the laser-pulse parameters q^+ and τ. It can be seen from the curves shown that, at any value of the quantity q^+, the fraction of reflected energy decreases, starting from some value of $\tau(q^+)$. This circumstance is connected with the growth in time of the optical thickness, $\int k\,dz$ ($k=2\omega\beta_2/c$), of the plasma layer and, correspondingly, with the decrease of the reflection coefficient $|V(t)|^2$. It may be assumed that in the case of spherical dispersion the quantity $R(q^+, t)$ attains a steady-state value, i.e., $R(q^+, t)\to R_{st}(q^+)$ for $t\to\infty$, as a result of the more rapid decrease of the density along the coordinate axis and the saturation of the quantity $\int k\,dr$.

In Fig. 4a, we show the profiles of the density, the temperature, and the energy-release (per unit mass) for $q^+=10^{14}$ W/cm^2 at the moment of time $t=1$ nsec. The energy-release curve has in the vicinity of the critical

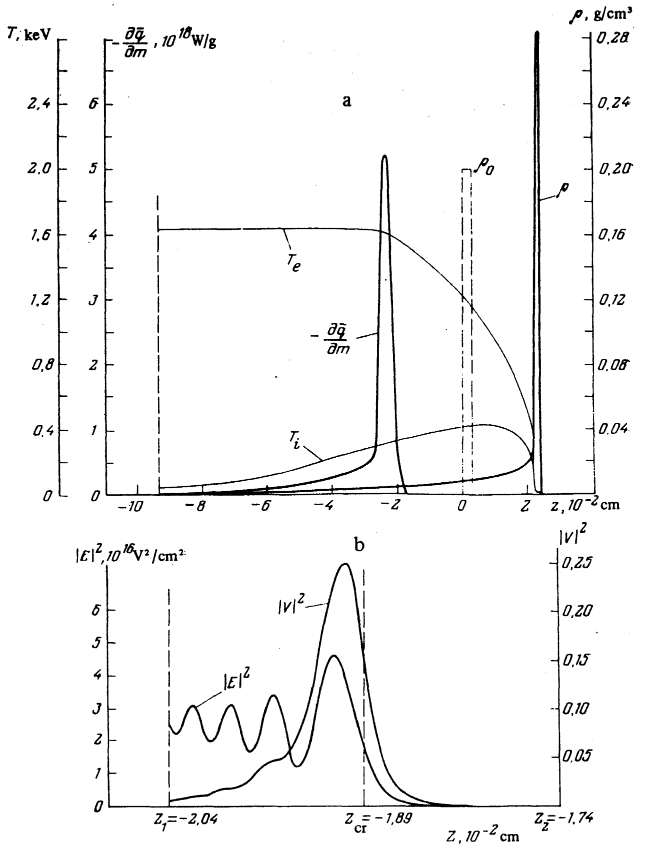

FIG. 4. The density, temperature, and energy-release profiles 1 nsec after the beginning of the action of a $q^+=10^{14}$ W/cm² constant pulse incident on a DT layer of thickness 30 μ, (a), and the structure of the field in the vicinity of the critical point at the same moment of time, (b).

point a strongly pronounced peak, which is connected with the anomalous-absorption mechanism corresponding to the frequency $\nu_1 \gg \nu_0$. Figure 4b shows the structure of the field $|E|^2$ and the quantity $|V|^2$ in the vicinity of the critical density. Let us note here the following circumstance. In the geometrical-optics approximation the quantity $|V|^2$ should decrease exponentially (without oscillations) toward the left plasma boundary. As follows from Fig. 4b, we can neglect the oscillations in $|V|^2$ near the point z_1 and go over to geometrical optics. In this case the quantity $|V|^2$ corresponds with a high degree of accuracy to the local coefficient of reflection with respect to energy.

To illustrate the dynamics of dispersion of a plasma layer, we present below the dependences of the brightening time, t^*, of a DT layer of thickness 30 μ on the flux q^+ for the cases of purely classical absorption (t^*_{cl}) and absorption with allowance for the anomalous processes (t^*_{an}):

q^+, W/cm²	10^{12}	10^{13}	10^{14}	10^{15}	10^{16}
t^*_{an}, nsec	52	9.1	1.86	0.818	0.403
t^*_{cl}, nsec	52	9.3	2.35	1.35	0.881

The time $t^*(q^+)$ was chosen from the condition of equality of the maximum density of the layer to the critical density.

Notice that in Fig. 2 we present values of the reflection coefficient integrated over the time interval t^*, which, as can be seen from the data presented above, depends on the magnitude of the incident flux. The dependence of the reflection coefficient on the magnitude of the flux for a fixed pulse duration can be found in Fig. 3, from which it can be seen that the reflection coefficient can increase, as well as decrease, with increasing flux.

We carried out a computation of the variant that has been realized in experiment[11]: on a solid polyethylene target impinges a flux of given shape and duration. Figure 5 shows the incident and reflected pulses, as well as the time dependence of the reflection coefficient (the values of $|V|^2$ in a vacuum), obtained in the computation. The experimentally observed fraction of reflected energy constitutes 3%, while the value obtained in the computation is 3.33%. It should be noted that the interpretation of the experiments with plane targets in the framework of the one-dimensional model has only a qualitative character. Indeed, the dimensions of the focal spot on the target are usually substantially smaller than the dimensions of the developing plasma "flare"; therefore, the problem should be two-dimensional (under the condition of axial symmetry). However, if in an experiment a focusing optical system with a small transmission value is used, then it can be assumed that the laser radiation interacts with some part of the "flare" where the curvature is insignificant. In Fig. 6 we show the corresponding profiles at the moment of time $t=4$ nsec after the pulse begins to act.

The influence of the contrast of the laser pulse on the magnitude of the reflected energy was also determined. First, a constant pulse ($q^+=10^{14}$ W/cm², $\tau=1.7$ nsec) incident on a layer of DT mixture was considered. The fraction of reflected energy in this case was 4.9% (see the curve corresponding to $q^+=10^{14}$ W/cm², $\tau=1.7$ nsec in Fig. 3). Then the case when this pulse had a forepulse linearly growing from zero to the value 10^{11} W/cm² in 3 nsec (without a time lag) was considered. The contrast in this case was $\sim 10^{-3}$. The fraction of reflected energy was 1.1%. In spite of the fact that the forepulse contained a negligible part of the energy, as compared to the main pulse, the fraction of reflected energy decreased to a value several times less than its former value.

It should be noted that in our calculations we used the

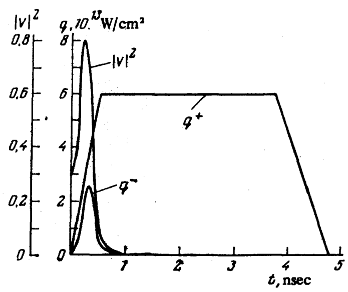

FIG. 5. The time dependences of the incident and reflected fluxes, as well as of the instantaneous reflection coefficient in the case of a solid polyethylene target.

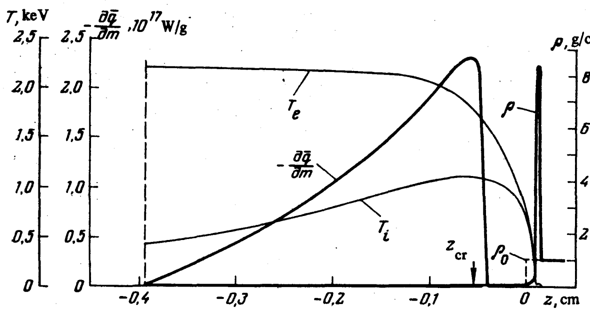

FIG. 6. The density, temperature, and energy-release profiles at 4 nsec (for a $(CH_2)_n$ target).

classical expression for the heat flux transportable by the electrons:

$$\Phi=-\varkappa_0 T_e^{5/2}\frac{\partial T_e}{\partial z}.$$

In[12] it is shown that the expression in question is valid when the condition $\Phi \leqslant \delta\Phi_{max}$, where

$$\Phi_{max}=n_e T_e (T_e/m_e)^{1/2}, \quad 0.03\leqslant\delta\leqslant 0.1.$$

is fulfilled. Estimates of the quantity δ for the temperature and density profiles shown in Fig. 4 yield ($\varkappa_0 = 1.93\times10^{19}$ g-cm-sec^{-3} – keV$^{-7/2}$) at the point $z = 2\times10^{-2}$ cm

$$T_e=0.5 \text{ keV} \quad |\partial T_e/\partial z|=40.6 \text{ keV-cm}^{-1} \quad n_e=4.7\cdot10^{21} \text{ cm}^{-3} \quad \Phi/\Phi_{max}=0.04;$$

at the point $z = 0$

$$T_e=1.23 \text{ keV} \quad |\partial T_e/\partial z|=21.4 \text{ keV-cm}^{-1} \; n_e=1.77\cdot10^{21} \text{ cm}^{-3} \quad \Phi/\Phi_{max}=0.1$$

and at the point $z = -1.9\times10^{-2}$ cm

$$T_e=1.56 \text{ keV} \quad |\partial T_e/\partial z|=14.4 \text{ keV-cm}^{-1} \quad n_e=10^{21} \text{ cm}^{-3} \quad \Phi/\Phi_{max}=0.2.$$

It follows from the obtained estimates that the use of the classical expression for Φ is in the present case physically correct.

The effective collision rates given in the table were obtained in the approximation of a weakly-turbulent, homogeneous plasma.[9] The approach developed in the present paper also allows us to use other possible effective collision rates, in particular, those obtained from numerical experiments. The inhomogeneity of the plasma can affect only the magnitude of the threshold for excitation of the dominant—in the cases considered in the present paper—instability, which consists in the decay of the light wave into a plasmon and ionic sound (of frequency ν_1). In order for the plasma inhomogeneity to significantly affect the magnitude of the instability-excitation threshold, the condition

$$H\ll l_e \ln(\omega/\nu_e),$$

where H is the characteristic inhomogeneity dimension, l_e is the electron mean free path, ω is the radiation frequency, and ν_e is the electron collision rate, should be fulfilled (see[9], p. 179). Let us verify this condition, using as an example the plasma whose parameters are given in Fig. 4: $H\approx10^{-1}$ cm, $v_s=(T_e/m_e)^{1/2}\approx1.6\times10^9$ cm-sec^{-1}, $\nu_{ei}\approx5\times10^{11}$ sec^{-1}, $l_e\approx3\times10^{-3}$ cm, and $l_e\ln(\omega/\nu_{ei})\approx2.4\times10^{-2}$ cm. These estimates show that the influence of the inhomogeneity in the present case is insignificant.

Thus, it can be seen from the computations carried out that, depending on the conditions (the pulse length, the presence of a forepulse), the fraction of reflected energy can be significantly different. These results can be used to explain experiments performed on different facilities (see the review chart in[13]). It should, of course, be noted that, besides the indicated effects, there exist a number of physical phenomena (e.g., refraction in the case of oblique incidence of the light) that must be taken into account in the explanation of the experimental data connected with the measurement of the light energy deposited in the plasma. Nevertheless, the results obtained in the present paper indicate that under certain conditions the magnitude of the laser-pulse energy deposited in the target plasma can be large enough to guarantee the requisite value of the hydrodynamic efficiency.

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