$\varepsilon^2 = \Delta^2 + s^2p^2$, and the envelope function satisfies the Klein-Gordon equation. The surface levels are given by the formula

$$E_k = -(\Delta^2 + s^2p_\perp^2)^{1/2}\left[1 + \frac{\Delta\alpha^2}{s^2k^2\sqrt{\Delta^2 + s^2p_\perp^2}}\right]^{-1/2},$$

where $p_\perp$ is the projection of the momentum on the surface of the crystal.

Thus, the system of levels lies near the ceiling of the lower band. The latter result can apparently not be applied directly to semiconductors with a narrow forbidden band. Indeed, in this case the lower band is the filled valence band of the semiconductor, and the problem assumes essentially a many-particle character. Its analysis lies beyond the scope of the present communication.

The following remark should be made concerning the calculations. They are valid if the frequency ω_e of electron motion on the surface levels is much smaller than the reciprocal relaxation time of the medium, a relaxation responsible for the establishment of κ. In this case the medium responds to the instantaneous position of the electron, and not to the averaged distribution of the charge with density $|\Psi|^2$. This is precisely why it is possible to use the usual expression for the potential of the electrostatic image. The parameter $(4\kappa)^2$ which enters in the theory is sufficiently large for both ionic (GaAs, InSb) and nonpolar (Ge, Si) semiconductors, so that ω_e turns out to be smaller than the frequency of the optical phonon, and certainly smaller than the frequencies of the bound electrons of the semiconductor.

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PLASMA HEATING BY ULTRASHORT LASER PULSES IN THE PROCESS OF ELECTRONIC HEAT CONDUCTION

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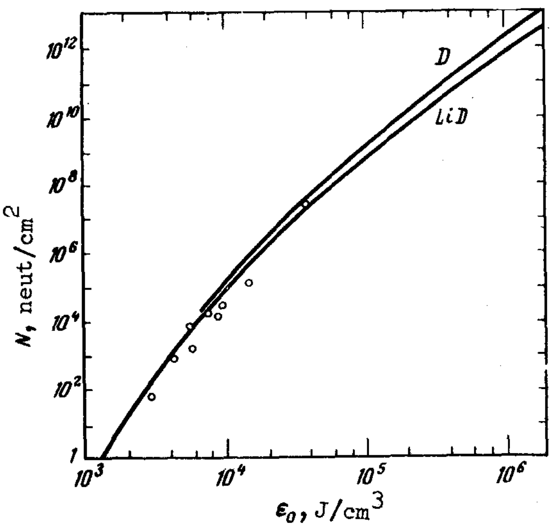
Plasma heating by focusing powerful ultrashort (10^{-11} - 10^{-12} sec) laser pulses on a solid target [1] may be accompanied by propagation of an electronic thermal wave in the interior of the target [2]. Let us consider the heating of ions in such a plasma at a laser-radiation absorbed-energy density $\varepsilon > \varepsilon_k \simeq 6 \times 10^3$ J/cm^2, when heat conduction makes possible an appreciable increase of the number of heated particles. The duration τ of the ultrashort pulse turns out in this case to be shorter than the characteristic plasma times, such as the time τ_{ei} of equalization of the electron and ion temperatures, the time τ_{ac} of propagation of a rarefaction wave through the heated layer, and the time τ_m of propagation of the heat-conduction wave (calculations show that the plasma radiation has little effect on the ion heating). For example, for LiD at $\varepsilon = 10^5$ J/cm^2 we have $\tau_{ei} \sim \tau_{ac} \sim \tau_m \sim 0.5 \times 10^{-10}$ sec. Under the indicated condition, the time of spreading of the layer (the time of expansion to twice the width) is $\tau_{gd} \sim 2 \times 10^{-10}$ sec. In addition, the thermal flux $S_e = \chi_e$ grad T_e does not depend on the electron density. It is therefore possible to disregard the presence of a "tail" of plasma with variable density near the target, and assume a model of instantaneous release of heat in the electrons in an infinitesimally thin surface layer of the target.

At the temperatures considered by us the plasma is fully ionized. The profile of the electronic thermal wave is given by $T_e(t)\,[1 - x^2/x_f^2]^{2/5}$, where x is the running coordinate reckoned from the surface of the target, and x_f is the coordinate of the wave front [3]. The ion temperature is taken in the same form, since their profiles coincide at the instant of temperature equalization, and during the earlier heating stage, in the case of a strong temperature difference, the concrete form of $T_i(x, t)$ has little effect on the results.



Yield of neutrons from a plasma vs. the absorbed laser-pulse energy. Solid curves - results of analytic calculations for solid deuterium and lithium deuteride. Points - results of numerical calculations [5] for solid deuterium at a laser pulse duration 6×10^{-12} sec.

After the termination of the laser pulse, a plane electronic thermal wave propagates in the interior of the medium, and heat is simultaneously transferred from the electrons to the ions as a result of the collisions; the plasma front spreads because of the large internal pressure. The heat-conduction wave front first leads the rarefaction wave, owing to the strong temperature dependence of S_e. The heat from the initially heated layer has time to propagate in the interior of the target before this layer has time to spread, and its energy is converted into translational-motion energy. It follows from the calculations that by the time the temperatures become equalized and the ion temperature reaches its maximum value T_{im}, the rarefaction wave propagates up to $x_{ac} = (2/3)x_f$. However, the layer thickness at which the density decreases by more than a factor of two as a result of the spreading amounts to approximately $x_{ac}/3$. Therefore the effect of spreading can be neglected when T_{im} is calculated. The time τ_m during which T_{im} remains relatively unchanged, to the contrary, is completely determined by the spreading of the layer of thickness x_m, which is heated by that instant.

The heat-conduction, energy-conservation, and electron-ion relaxation equations, written for the temperatures

$$T = \frac{1}{x_f} \int_0^{x_f} T(t) \left[1 - \left(\frac{x^2}{x_{\Phi}^2}\right)\right]^{2/5} dx,$$

averaged over the profile, take the form

$$x_f \frac{dx_f}{dt} = a T_e^{5/2}, \qquad (1)$$

$$T_e + \frac{1}{z} T_i = \frac{2}{3} \frac{\epsilon}{n_e} \quad (2)$$

$$\frac{dT_i}{dt} = \kappa \frac{T_e - T_i}{T_e^{3/2}}, \quad (3)$$

where

$$a = \frac{0{,}86\,\delta(z)}{m_e^{1/4} e^4 n_e \ln \Lambda}, \qquad \kappa = \frac{8\sqrt{2\pi}\, m_e^{1/2} e^4 z^2 n_e \ln \Lambda}{3 m_i},$$

$\delta(z) = 1$ when $z = 1$ [4], z and $\tilde{z}$ are the mean-squared and mean charges of the ions, and m_i is the average ion mass. The solution of (1) - (3) yields

$$T_{im} = 0{,}77 \cdot 10^9\, e^{4/3} \left(\frac{m_e}{m_i}\right)^{1/6} (\ln \Lambda)^{1/3} [\delta(z)]^{-1/6} z^{1/3} \xi_m(\tilde{z}) \epsilon^{1/3} \text{ (keV)} \quad (4)$$

$$\xi_m(1) = 0{,}44\,.$$

Here

$$x_m = \frac{2}{3} \frac{\epsilon}{n_e T_{im}\left(1 + \frac{1}{\tilde{z}}\right)}, \quad (5)$$

$$r_m = 1{,}13 \left[\frac{m_i}{3 T_{im}(1 + \tilde{z})}\right]^{1/2} x_m \,. \quad (6)$$

The growth of the ion temperature with increasing energy absorption indicates that the heat-conduction heating is effective. Expressions (4) - (6) make it possible to determine the neutron yield from the plasma (see the figure). In the case of LiD, a calculation was performed also in the region where there is no heat conduction. The results [1] correspond to $\varepsilon \simeq 5 \times 10^3$ J/cm^2, i.e., where the heat conduction still does not exert a noticeable influence.

A recent communication reports numerical calculations of ion heating in solid deuterium with the aid of subnanosecond laser pulses [5]. The values of the relative neutron yields obtained in these calculations are in good agreement with the results corresponding to the analytic expressions (4) - (6) (see the figure).

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