Let us compare the effectivenesses of the two mechanisms. Assuming $\beta \sim e/a_0^2$ and $\Lambda \sim (\hbar^2/ma_0^2)$ (a_0 is the lattice constant), we obtain an order-of-magnitude estimate of the ratio

$$\left|\frac{\delta\epsilon_{el}}{\delta\epsilon_{def}}\right| \sim \frac{m}{m^*}\frac{\lambda_T}{r_D}\sqrt{\frac{E_g-\hbar\omega}{T}}\left(\lambda_T = \frac{\hbar}{m^* v_T}\right). \tag{7}$$

Near the boundary frequency ($E_g - \hbar\omega \sim T$), the effectivenesses turn out to be of the same order. In particular, for CdS this ratio is of the order of 0.6. According to estimates for CdS, the acousto-optical effectiveness near the edge of the band is equal to $M_2 \simeq 10^{-16}$ sec^{2}g^{-1} and exceeds by one order of magnitude the values of M_2 obtained in experiment far from the edge of the band [6]. The estimates have a qualitative character, since in the anisotropic medium the contribution made to the scattering by any particular mechanism depends strongly on the crystal orientation.

In conclusion we note that the acousto-optical effectiveness of the electronic mechanism depends strongly on the concentration of the conduction electrons. In photosemiconductors (CdS, CdSe), application of infrared illumination makes it possible to control the conduction-electron density [7], and this, in turn, makes it possible to control effectively the scattering and diffraction of light in such crystals.

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ROLE OF STIMULATED COMPTON SCATTERING IN THE INTERACTION OF LASER RADIATION WITH A SUPERDENSE PLASMA

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Recently, much hope has been placed in the possibility of obtaining a controlled thermonuclear reaction by fast adiabatic compression of thermonuclear matter heated with a laser [1]. We wish to call attention to one circumstance that must be taken into account both in the theoretical description of different models of fast heating of a superdense plasma by laser radiation and in their practical realization. We have in mind the possibility of strong reflection of the laser radiation from the heated plasma due to the stimulated Compton scattering of light by the plasma electrons. As is well known, in such a process the gasdynamic expansion of the hot plasma leads to formation of a plasma corona around the dense part, in which the plasma, with varying density gradient, moves in a direction opposite to the laser beam. Our earlier experimental and theoretical investigations [2] have shown that in the presence of a

Fig. 1. Total reflection coefficient $R = I_2(\ell)/I_{10}$ vs. initial reflection $R_0 = I_2(0)/I_1(0)$ for $\ell = 0.2$ cm (1), 0.1 cm (2), 0.05 cm (3), and 0.01 cm (4) at $I_{10} = 10^{16}$ W/cm^2.
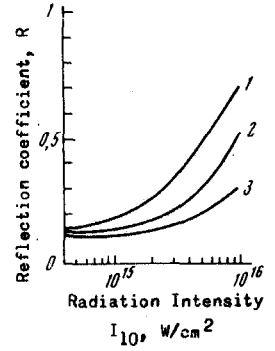
plasma moving opposite to the laser beam, the fraction of the reflected light energy can reach a noticeable value. Under certain conditions, the reflection may be total. In the presence of a density gradient in the region where the plasma frequency is equal to the laser-radiation frequency, the reflection coefficient is always finite. The resultant opposing wave interacts during the stimulated Compton scattering with the intense incident wave and is noticeably amplified. At sufficiently high radiation fluxes, the energy of the incident wave is effectively converted into reflected-wave energy.

The equations describing the amplification of the weaker pulse in the presence of a moving plasma and intense radiation were obtained in [2] in the form

$$dI_1(z)/dz = dI_2(z)/dz = \beta I_1(z) I_2(z) . \tag{1}$$

The positive z direction is defined as the direction of plasma motion towards a beam of intensity I_1 and coincides with the direction of propagation of the reflected radiation with intensity I_2, while β is defined by the equation

$$\beta = \frac{c r_0^2}{\pi \nu^3} \frac{v_0}{kT_e} n_e \left(\frac{m}{2\pi kT_e} \right)^{1/2} \exp\left(- \frac{mv_0^2}{2kT_e} \right) , \tag{2}$$

where c is the speed of light, r_0 is the classical radius of the electron, ν is the laser radiation frequency, m is the electron mass, k is Boltzmann's constant, v_0 is the plasma velocity, n_e is the plasma density, and T_e is the plasma electron temperature. In the derivation of these equations it was assumed that $\bar{v}/c \gg \Delta\nu/\nu$ ($\bar{v}$ is the average thermal velocity of the electrons and $\Delta\nu$ is the width of the emission spectrum). Solution of (1) leads to the following relation, from which we can determine the reflection coefficient $R = I_2(\ell)/I_{10}$ of laser radiation from a plasma layer of thickness ℓ:

$$R = R_0 \exp \beta I_{10} \ell (1 - R), \tag{3}$$

where I_{10} is the intensity of the laser radiation entering the plasma layer (at $z = \ell$), and $R_0 = I_2(0)/I_1(0)$ is the initial reflection coefficient (at $z = 0$). Using the average parameters obtained in the calculation of the adiabatic compression model [1], namely $kT_e = 10$ keV, $n_e \simeq (0.5 - 1) \times 10^{21}$ cm^{-3} and $v_0 = 10^8$ cm/sec at the neodymium-laser frequency $\nu = 2.83 \times 10^{14}$ sec^{-1}, we obtain from (2) $\beta = 31.4 \times 10^{-16}$ cm/W. The solution of (3) is shown in Fig. 1 in the form of plots of the total reflection coefficient R against the initial reflection coefficient R_0, for four values of the parameter ℓ (0.2, 0.1, 0.05, and 0.01 cm) at an intensity $I_{10} = 10^{16}$ W/cm^2 [1]. It is seen from the curves of Fig. 1 that at the given parameter values the effect can lead to an appreciable increase of the laser-radiation reflection from the plasma. Actually, one can use as the parameter not the quantity ℓ, but the product $\beta\ell$, and the plots can be used to determine the dependence of the reflection on any of the
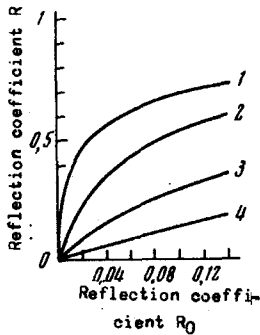
Fig. 2. Total reflection coefficient $R = I_2(\ell)/I_{10}$ vs. the intensity of the incident radiation I_{10} at $R_0 = 0.1$ and $\ell = 0.2$ cm (1), 0.1 cm (2), and 0.05 cm (3).

quantities ν, ν_0, n_e, and kT_e (see formula (2)). It should be noted that although according to (2) β depends very strongly on the frequency ($\beta \sim \nu^{-3}$), it can be assumed approximately that the total reflection coefficient depends much less on the frequency in the spherical-heating model, owing to the frequency dependences of the parameters ℓ and I_{10}.

From the plots of the reflection coefficient against the incident radiation intensity I_{10}, shown in Fig. 2 for $\ell = 0.2$, 0.1, and 0.05 cm and for $R_0 = 0.1$, it follows that the reflection increase due to the mechanism under consideration becomes noticeable at fluxes $I_{10} \simeq 10^{15}$ W/cm^2.

An exact calculation of the reflected energy should include, naturally, integration with respect to the spatial distribution of the temperature and of the plasma expansion rate and density, as well as with respect to the distribution of the intensity I_1. It should also be borne in mind that the electric field plays an important role in stimulated Compton scattering. At densities such that the plasma frequency is equal to the laser frequency, the effective field increases strongly, and this enhances greatly the effectiveness of the processes of stimulated Compton scattering, and can lower the flux at which the increase of the reflection becomes appreciable.

It is physically clear that this effect is due to interaction of single electrons with photons, and is consequently practically inertialess.

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