intensity of the quasistationary noise to the shot noise. To this end we put

$$\epsilon_1 = \epsilon'(\omega) - i\frac{4\pi\sigma(\omega)}{\omega} \text{ and } |\epsilon_1| >> 1, \tag{4}$$

where σ is the conductivity, and obtain

$$\gamma(\omega) = \frac{I\sqrt{eE_0}}{\pi k T\epsilon'}\frac{\tau}{1+\omega^2\tau^2}, \tag{5}$$

where $\tau = \epsilon'/4\pi\sigma$.

We see that the quasistationary noise in semiconductors has a characteristic frequency dependence. Moreover, if $\epsilon'(\omega)$ and $\sigma(\omega)$ depend little on the frequency, then this dependence is typical of noise in semiconductors; for example, it is characteristic of generation-recombination noise [4]. In semiconductors with not too high a conductivity, one can observe here a quasistationary noise which exceeds greatly the shot noise in intensity, at relatively small emission currents. Thus, at $\sigma \sim 10^{11}$ sec^{-1} ($\rho \sim 10$ ohm-cm), $E_0 = 10^2$ V/cm, $T \sim 10^3$ °K, I = 1 mA, $\gamma(0) \sim 70$. Whereas in semiconductors the main contribution to the frequency dependence of $\gamma(\omega)$ is made by the first term of (1), in metals the main contribution is made by the second term, while the first gives rise to an increment that is independent of the frequency [5].

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PARAMETRIC BACK SCATTERING OF A LINEAR ELECTROMAGNETIC WAVE IN A PLASMA

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Possible mechanisms of energy absorption from a powerful laser beam by the plasma corona of a contracting deuterium drop [4, 5] are connected with parametric instabilities [1 - 3]. Two parametric processes lead to the conversion of the energy of the incident electromagnetic wave into the energy of plasma oscillations, namely the "decay instabilities" - the pump wave (photon) "plasmon" + "phonon" (ion sound) and the pump wave "plasmon" + "plasmon." The first of these processes can occur in a certain layer of the corona, where the frequency of the incident electromagnetic wave ω_0 is approximately equal to the electron plasma frequency

$$\omega_0 \cong \omega_p = \sqrt{\frac{4\pi n_c e^2}{m}}.$$

The second process can occur at approximately one-quarter of the critical density, at $n = n_c/4$.

However, there are possible parametric instabilities, as a result of which the energy of the pump wave is transformed mainly into scattered electromagnetic radiation. There are "decays" of two types: "photon" → "photon" + "plasmon" and "photon" → "photon" + "phonon." The kinematic conditions for parametric resonance, $\omega_0 = \omega_1 + \omega_2$ and $\vec{k}_0 = \vec{k}_1 + \vec{k}_2$ do not forbid such processes also in the outer part of the corona, where the density $n \ll n_c/4$. These processes can present a serious danger if the incident electromagnetic wave has time to be scattered before reaching the region where $n = n_c/4$. The question is whether the scattered electromagnetic waves have time to become parametrically amplified as they start from the initial thermal-noise level and propagate in the inhomogeneous plasma corona. Let $\gamma(r)$ be the local increment of the parametric instability in the geometrical-optics approximation. Then the intensity of the scattered "photon" increases in accordance with the law

$$E_1^2 \sim \exp\left\{ 2 \int \gamma(r) \frac{dr}{(\partial\omega/\partial k_r)_1} \right\}. \tag{1}$$

Here $(\partial\omega/\partial k_r)_1$ is the group velocity of the "photon" in the direction of r (along the density gradient). The expression for the increment of the decay instabilities for the case of a homogeneous plasma is [1, 6]

$$\gamma = [\,|V_{0,1,2}|^2 |C_0|^2 - \delta^2/4]^{1/2}\,, \tag{2}$$

where $V_{0,1,2}$ is the matrix element of the process $0 \to 1 + 2$, $\delta = \omega_0 - \omega_1 - \omega_2$ is the frequency detuning of the parametric resonance, and $C_0 = E_0/\sqrt{8\pi\omega_0}$. The matrix elements of the processes investigated by us were calculated earlier [1, 7]. The maximum gain (1) is obtained when the photon is scattered through an angle θ close to 90°. The denominator of the integrand in (1) is then minimal. The dependence of the group velocity $\partial\omega/\partial k_r$ on r is determined from the condition

$$\omega^2 = \omega_p^2(r) + c^2k^2 = \text{const}\,. \tag{3}$$

The main contribution to the parametric-resonance frequency detuning is made by the inhomogeneous Doppler effect when the corona expands with supersonic velocity $U(r)$ (for the process in which the "phonon" takes part) and by the inhomogeneity of the corona density (for the process with the "plasmon"). The gain for the first of these processes is appreciably larger (see the table). The corresponding expression for the frequency detuning is

$$\delta(r) = \frac{\omega_0}{c}\frac{dU(r_0)}{dr_0}(r - r_0), \qquad 1 > \frac{\gamma_d}{k_0 U} > \frac{\omega_p^2}{\omega_0^2}\,. \tag{4}$$

In the general case, the process of parametric back scattering with allowance for the nonlinearity of the secondary waves is quite complicated. However, it is possible to construct an analytic model in a quasilinear approximation, when the interaction between the modes can be neglected, but their reaction on the pump-wave intensity is taken into account

$$\frac{1}{r^2}\frac{d}{dr} r^2 c \frac{\mathbf{E}_0^2}{8\pi} = 2 \int \gamma(r, \mathbf{k}_1)\, \mathbf{E}_{\mathbf{k}_1}^2 \frac{d^3\mathbf{k}_1}{(2\pi)^3}\,, \tag{5}$$

Process	Instability increment	Gain for scattering through 90°
Decay into "photon" + "phonon" $T_e \gg T_i$	$\gamma_d^2 = \frac{\omega_p^4}{8\omega_0^2}\left(\frac{2T_e}{Mc^2}(1-\cos\theta)\right)^{1/2}\frac{E_0^2}{8\pi n T_e}$	$\nu = 12{,}4\frac{\gamma_d^{3/2}}{\omega_p\sqrt{\bar{U}'}}(k_0 L)^{1/2}$
Stimulated scattering by ions $T_e = T_i$	$\gamma = \frac{\omega_p^4}{4\omega_0^3}\frac{E_0^2}{8\pi n T_e}\,\mathrm{Im}\frac{1}{(k_0-k_1)^2\lambda_0^2\epsilon(\omega_0-\omega_1)}$	$\nu \sim \frac{\omega_p^3}{\omega_0^3}k_0 L^{1/2}\left(\frac{T_e}{MU'^2}\right)^{1/4}\frac{E_0^2}{8\pi n T_e}$
Modified decay $\gamma_d > k_0 v_s$	$\gamma = \frac{\sqrt{3}}{2}\left[\frac{\omega_p^4}{2\omega_0}\frac{T_e}{Mc^2}(1-\cos\theta)\frac{E_0^2}{8\pi n T_e}\right]^{1/3}$	$\nu = 4B\left(\frac{1}{2}\ \frac{a-1}{3a}\right)\frac{\gamma}{\omega_p}k_0 L,\ n \sim r^{-a}$
Decay into "photon" + "plasmon"	$\gamma_d^2 = \frac{\omega_p^3}{4\omega_0}\frac{E_0^2}{8\pi n mc^2}(1-\cos\theta)$	$\nu = 17{,}6\left(\frac{\gamma_d}{\omega_p}\right)^{3/2}k_0 L$

where the amplitude E_{k_1} is obtained with the aid of (1) by using the level of the thermal noise as the initial amplitude. The main contribution to the integral of (5) is made by the secondary wave which is maximally amplified (we note that these are different waves for different r).

The solution of (5) is of the form[1])

$$\nu^{-7/6}e^{\nu} = \Lambda L\, d\ln r^2 E_0^2/dr\ ,$$

$$\Lambda = 10^2 n\left(\frac{c}{\omega_p}\right)^3\left(\frac{\omega_0}{\omega_p}\right)^{7/3}(k_0 L)^{-2/3}\left(\frac{\bar{U}c^2}{v_s^3}\right)^{1/3}\ ,$$

$$L^{-1} = d\ln n/dr \ll \frac{\gamma}{c},\quad \bar{U} = L\,dU/dr,\quad v_s = \sqrt{\frac{T_e}{M}}\ . \tag{6}$$

The pump wave is transformed almost completely into scattered electromagnetic radiation, and the corresponding incident-energy flux decreases rapidly inside the corona. We point out that the model is valid for not too large intensities of the incident radiation, when the weak-turbulence approximation is still valid. At high light intensity, when $\gamma > k_0 v_s$, formula (2) for the increment γ is modified: $\gamma = [2^{3/2}k_0(1-\cos\theta)^{1/2}v_s\gamma_d^2]^{1/3}$ (this case is designated "modified decay" in the table). The qualitative picture of the back scattering for this case remains the same as before.

At very large light intensities, however, at any rate when $E^2 > 8\pi n_c T$, the picture of the parametric reflection should be completely altered, since the light pressure causes a sharp decrease of the plasma density at $n < n_c$ (inside a nonlinear skin layer of sorts [8, 9]). An entirely different approach to the parametric instabilities and consequently to absorption is required in the interior of such a skin layer. We note that the effective absorption of the light in the skin layer (when the absorption coefficient is greater than v_{Te}/c) should

[1])The incident radiation is assumed here to be unpolarized, so that the gain ν in this formula has half the value listed in the table for $\vec{E} = \vec{E}_0\cos(\omega_0 t - k_0 r)$.

lead to overheating and to the appearance of the electronic thermal-conductivity instability [10].

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HEATING AND CONTAINMENT OF PLASMA IN CROSSED LIGHT BEAMS

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Heating of plasma by pulsed-laser radiation for the purpose of obtaining a thermonuclear reaction is the subject of a number of theoretical papers (see, e.g., [1 - 4]). These papers deal with conditions in which the heating time does not exceed the time of the hydrodynamic expansion of the plasma. In the present communication we propose another method of plasma heating by pulsed-laser radiation. In this method, the time of heating can be determined by the duration τ_p of the laser pulse, and by the same token can be much larger than the corresponding time of the hydrodynamic expansion. In addition, unlike the conditions considered in the cited papers, in our method the area of the skin layer (in which the direct conversion of the light energy into thermal energy or into energy of collective plasma oscillations takes place) can be much larger then the surface area bounding the entire volume of the heated plasma.

For such heating, it is necessary to place the substance in the region of intersection of two light beams obtained, for example, from one laser and inclined to each other at a certain angle α. We consider below for concreteness focused beams (with individual focal-region diameters d_{f_1} and d_{f_2}, intersecting in the vicinity of these focal regions. By direct calculation we can verify that under asymmetrical conditions, for example (i) when $d_{f_1} \neq d_{f_2}$ or (ii) when the centers of the focal regions do not coincide and the beams in question have no symmetry plane, etc., three-dimensional "microregions" are produced as a result of the interference of the fields of the crossed beams. At the nodal points of each of the microregions, the values of the optical energy density (and consequently of the electromagnetic-field pressure) are small in comparison with the corresponding values on the boundary. One of the dimensions (Δx) of the microregion is obviously $\Delta x \simeq \lambda/2 \sin(\alpha/2)$ (λ is the laser emission wavelength), i.e., is determined by the distance between the interference bands of two plane waves; the two others (Δy and Δz) can be determined (i) by the