# Amplification of reflected light in an expanding laser plasma

L. M. Gorbunov

*P. N. Lebedev Physics Institute, USSR Academy of Sciences*
(Submitted 7 August 1980)
Zh. Eksp. Teor. Fiz. **81**, 146–150 (July 1981)

The amplification of light reflected from a dense plasma as a result of stimulated scattering in the rarefied expanding plasma corona is considered. It is shown that the radiation intensity that reaches the dense layers of the plasma depends on the scale over which the rate of plasma flow changes, and does not depend on the incident radiation intensity if the latter exceeds a certain value.

PACS numbers: 52.25.Ps

When laser radiation acts on a target, the incident light wave and that reflected from the dense plasma produce a standing electromagnetic wave in the expanding rarefied plasma corona. The ponderomotive forces that act in this wave produce in the plasma density perturbations that are particularly pronounced in the resonance region, where the rate of plasma expansion is of the order of the speed of sound, and the Mach number is close to unity. Scattering by these perturbations increases the reflection of the light from the corona. This effect was first discussed in the papers of Begg and Cairns[1] and of Randall *et al.*,[2] who used for the plasma corona highly simplified models, incapable of revealing the main laws that govern the phenomenon.

The present paper the amplification of reflected light is analyzed using a more realistic model. The spatial variation of the density perturbations and of the intensities of the incident and reflected light are investigated and the reflection coefficient is determined. It is shown that the characteristic scale of the change of the flow velocity is a most important factor in the effectiveness of the reflected-light amplification.

1. The state, unperturbed by an electromagnetic field, of an inhomogeneous plasma corona will be characterized by the electron density $N(x)$, by the plasma flow velocity $U(x)$, and by a constant temperature $T$. The time-independent small deviations from this state, which are connected with the action of the electromagnetic field, are determined by the equations[3]

$$\frac{d}{dx}(nU+vN)=0;$$
$$\frac{d}{dx}(vU)=-v_s^2\frac{d}{dx}\left(\frac{n}{N}\right)-\frac{ze^2}{4mm_i\omega_0^2}\frac{d}{dx}|E|^2-\nu v. \tag{1}$$

where $n$ and $v$ are the preturbations of the density and of the flow velocity, $v_s=(zT/m_i)^{1/2}$ is the speed of sound; $z$ and $m_i$ are the charge number and the mass of the ions; $\nu$ is the effective collision frequency; $\omega_0$ and $E$ are the frequency and intensity of the electric field.

At normal incidence of the radiation on the plasa, the value of $E$ is determined from the equation

$$\frac{d^2E}{dx^2}+k_0^2(x)E=\frac{n\omega_p^2(x)}{Nc^2}E, \tag{2}$$

where

$$k_0^2=\frac{\omega_0^2}{c^2}\left(1-\frac{N(x)}{N_c}\right),\quad N_c=\frac{m\omega_0^2}{4\pi e^2},\quad \omega_p^2=\frac{4\pi e^2N}{m}.$$

We seek the solution of (2) in the form

$$E=E_0(x)\exp\left[-i\int^x dx'k_0(x')\right]+E_1(x)\exp\left[i\int^x dx'k_0(x')\right],$$

where $E_0$ and $E_1$ are the amplitudes of the waves incident on and scattered from the plasma. Using this expression and assuming that

$$n=n_0(x)\exp\left[2i\int^x dx'k_0(x')\right]+n_0^*(x)\exp\left[-2i\int^x dx'k_0(x')\right],$$

we obtain from Eqs. (1) in the zeroth geometric-optics approximation

$$n_0(x)=-\frac{z\omega_p^2(x)/\omega_0^2}{16\pi T[1-M^2(x)+i\nu M/2k_0v_s]}E_0^*E_1,$$

where $M(x)=U(x)/v_s$ is the Mach number.

In the first geometric-optics approximation we then get from Eq. (2) two equations for the intensities of the incident and reflected (scattered) waves[1)]:

$$dq_0/d\xi=\varkappa(\xi)q_0q_1,\quad dq_1/d\xi=\varkappa(\xi)q_0q_1, \tag{3}$$

where

$$\xi=\omega_0x/c,\quad q_{0,1}=c^2k_0|E_{0,1}|^2/8\pi\omega_0,\quad q_T=N_cTc,$$
$$\varkappa(\xi)=\frac{(N(\xi)/N_c)(\nu M/2k_0v_s)}{2q_T(1-N(\xi)/N_c)[(1-M^2)^2+(\nu M/2k_0v_s)^2]}. \tag{4}$$

We assume that at $\xi=\xi_0$ the plasma density vanishes and in this case the intensity of the incident wave is known and equals $I=q_0(\xi_0)$. We choose as the origin a point located close enough to the reflection point [where $N(\xi)=N_c$], but in which the geometric optics approximation is still valid. The ratio $(q_1(0)/q_0(0))=r$ at this point is assumed to be known. It follows then from the solution of Eqs. (3) that the total coefficient of reflection from the plasma corona $R=q_1(\xi_0)/I$ is given by

$$R=r\exp[I(1-R)G(\xi_0)], \tag{5}$$

where

$$G(\xi)=\int_0^\xi d\xi'\varkappa(\xi').$$

The functions $q_{0,1}$ then take the form

  

$$q_0=\frac{I(1-R)}{1-r\exp[I(1-R)G(\xi)]},\quad q_1=\frac{I(1-R)}{r^{-1}\exp[-I(1-R)G(\xi)]-1}. \tag{6}$$

The density perturbations are expressed interms of the functions (6):

$$|n|^2=N^2(\xi)\frac{q_0q_1}{4q_T^2(1-N/N_c)[(1-M)^2+(\nu M/2k_0v_s)^2]}.$$

2. The general properties of Eq. (5) and of the functions (6) do not depend on the concrete hydrodynamic characteristics of the corona. According to the definition (4) we have $\varkappa(\xi)>0$ and the function $G(\xi)$, as well as the functions $q_0$ and $q_1$, increase monotonically with increasing $\xi$. Consequently at $\xi=0$ they have minimal values and equal, according to Eq. (6),

$$q_{0,\,min}=q_0(0)=I(1-R)/(1-r),\quad q_{1,\,min}=q_1(0)=I(1-R)r/(1-r). \tag{7}$$

If the reflection is weak ($R<1$), the formula (5) is simplified and $R\approx r\exp[IG(\xi_0)]$. This relation is valid only for the values

$$I<I_0=G^{-1}(\xi_0)\ln(1/r). \tag{8}$$

If an inequality inverse to (8) is satisfied, then the reflection coefficient is close to unity, $R\approx 1-I_0/I$. Therefore at $I\sim I_0$ the exponential growth of the function $R(I)$ slows down (Fig. 1) and saturation of the scattering sets in. It follows then from (7) that $q_0(0)=I_0/(1-r)$ and the laser radiation reaching the dense plasma layer has a constant intensity independent of the intensity of the incident radiation. A similar result is known for the SMBS in liquids,[10] and was obtained also in a calculation of the attenuation, on account of SMBS, of the pump wave in an inhomogeneous laser plasma.[5]

If the point $\xi_1$ at which the equality $M(\xi_1)=1$ is satisfied is located in the rarefied-plasma region ($\xi_0>\xi_1>0$), then the reflection coefficient and the functions (6) can be approximately expressed in terms of the hydrodynamic characteristics of the corona. To this end it is necessary to expand the function $M(\xi)$ in (4) in a series in the vicinity of the point $\xi_1$. As a result we get

$$G(\xi)=G_0\left\{\operatorname{arctg}\frac{(\xi-\xi_1)}{a}+\operatorname{arctg}\frac{\xi_1}{a}\right\}, \tag{9}$$

where $a=\nu(\xi_1)/4k_0(\xi_1)v_sM'$ characteristics the width of the region of the resonant scattering,

$$G_0=\frac{1}{4q_T}\left(\frac{N(\xi_1)}{N_c}\right)\left[M'\left(1-\frac{N(\xi_1)}{N_c}\right)\right]^{-1},\quad M'=\left.\frac{dM}{d\xi}\right|_{\xi=\xi_1}.$$

If $a$ satisfies the conditions $\xi_1$, $(\xi_0-\xi_1)>a$, then $G(\xi_0)=\pi G_0$, and $I_0$ [see (8)] takes the form

$$I_0=\frac{4}{\pi}\ln\left(\frac{1}{r}\right)q_T\frac{M'(1-N(\xi_1)/N_c)}{N(\xi_1)/N_c}. \tag{10}$$

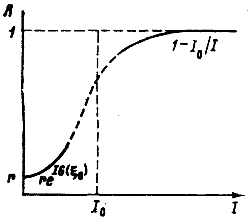


FIG. 1. Reflection coefficient vs. intensity of incident radiation.

It is seen from (10) that the radiation intensity at which the growth of the reflection coefficient slows down does not depend on the width of the resonance region, but is determined by the characteristic scale of the variation of the Mach number. If the reflection from the dense plasma is weak ($r\ll 1$), then $q_0(0)\approx I_0$ and expression (10) determines the same limiting laser-radiation intensity that reaches the dense plasma layers.

Figure 2 shows the functions $q_{0,1}(\xi)$ and $|n(\xi)|$ at two different values of the parameter $a$. It is seen that with decreasing effective collision frequency $\nu$ the region of resonant excitation of the density perturbations becomes narrower, and their amplitude increases. Neither the total change of $q_{0,1}$ nor the reflection coefficient $R$ is then dependent on $\nu$, and both are determined by the characteristic scale of the change of the plasma flow velocity.

4. We shall now make a few remarks concerning the conditions under which the results are valid. In the analysis of the amplification of the reflected light, no account was taken of its absorption via collisions of the electrons with the ions (collision frequency $\nu_{ei}$). This assumption is justified if the nonlinear interaction between the incident and reflected waves leads to a more rapid variation of their intensities than the collisions. This calls for satisfaction of the condition

$$\frac{\nu_{ei}}{\omega_0}<\frac{2k_0v_s}{\nu}\left(1-\frac{N}{N_c}\right)^{-1/2}\frac{q_0}{2q_T}. \tag{11}$$

Both for the density perturbations and for the electromagnetic field, we have confined ourselves to the first harmonics. The condition that the higher harmonics of the density perturbations be small at the point $\xi_1$ is the most stringent one and is of the form

$$(\nu/2k_0v_s)^2>4q_0/q_T. \tag{12}$$

In a rarefied plasma, the main mechanism of dissipation of the density perturbations is Landau damping. If the plasma is non-isothermal, then the damping is by electrons ($\nu/2k_0v_s\sim(m/m_i)^{1/2}\ll 1$) and the inequality (12) is satisfied only for relatively low intensities of the incident radiation ($q_0/q_T<m/mi$). On the other hand if the temperatures of the electrons and ions are close, then strong Landau damping by ions sets in, and $\nu/2k_0v_s\lesssim 1$. Under these conditions our analysis is valid for $q_0\lesssim q_T$ and in this case the inequality (11) is reliably satisfied in a laser plasma ($\nu_{ei}/\omega_0\ll 1$).

A remark must be made also concerning the plasma

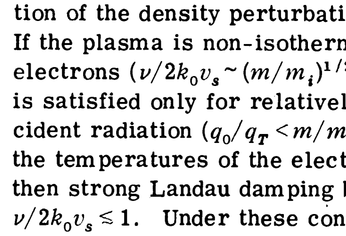


FIG. 2. The intensities of the incident ($q^0$) and reflected ($q_1$) light and the amplitude of the perturbed plasma density ($|n|$) vs. the coordinate $\xi$ for two values of the parameter $a$. The point $\xi_1$ is defined by the condition $M(\xi_1)=1$.

flow unperturbed by an electromagnetic wave. It is known that the one-dimensional isothermal plasma flow is not stationary. However, if the changes of the functions $N$ and $U$ take place over a characteristic time $\tau$ longer than the settling time of the density perturbations in an electromagnetic field, then the nonstationary character of the flow can be disregarded. The density-perturbation settling time, for sufficiently strong dissipation, is $\nu^{-1} \sim (2k_0 v_s)^{-1}$. The time of variation of the plasma parameters can be estimated, in order of magnitude, at $\tau \sim L/v_s$, where $L$ is the scale of variation of the functions $N$ and $U$. It follows therefore that the plasma flow is quasistationary if $k_0 L > 1$. Satisfaction of this inequality was already assumed in the geometric-optics method used by us.

The question of the validity of the linear approximation in stimulated scattering was discussed in Ref. 11, where stimulated scattering by electrons and ions in an isothermal plasma was considered and it was shown that under laser-plasma conditions appreciable scattering sets in only at large density perturbations ($n/N \sim 1$), which cannot be described by the linear approximation alone. In our case the principal scattering is by acoustic waves. We use here the linear theory and assume by the same token that $n/N \ll 1$. From the expression obtained for this quantity in the resonance region ($M = 1$) and under conditions of considerable scattering ($I \approx I_0$) it follows that the linear approximation is valid if the following inequality holds:

$$\frac{\nu}{2k_0 v_s} \gg \frac{2}{\pi} \ln\left(\frac{1}{r}\right) \frac{(1-N(\xi_1)/N_c)^{1/2}}{N(\xi_1)/N_c} \frac{c}{\omega_0} \frac{dM}{dx} \sim \frac{1}{k_0 L},$$

Under conditions when geometric optics is valid, this inequality is quite well satisfied.

5. Resonant excitation connected with reflection of an electromagnetic wave from a dense plasma was observed in a numerical simulation of laser plasma.[2] Attributed to the same effect are the results of experiments on the interaction of microwave radiation with a rarefied plasma.[6-8] In particular, in the experiment of Akiyama *et al.*,[8] the most intensive sources of density perturbation were observed in a region where the plasma flow velocity was close to that of sound.

An indirect indication of the onset of this effect is the saturation, at $I_0 \sim 10^{13}$ W/cm$^2$, in the reflection of the radiation of a $CO_2$ laser from a rarefied plasma at a level $R \sim 1$ (Ref. 9). If we use the plasma parameters indicated in Ref. 9, ($N/N_c \approx 0.18$; $L = 200$ $\mu$m; $T = 100$ eV), then we obtain from formula (10) $I_0 \approx 0.8 \times 10^{13}$ W/cm$^2$, which agrees well with the experimental data.

In conclusion, I wish to thank N. E. Andreev and V. P. Silin for a number of helpful remarks.

[1)] For a homogeneous medium, a system of equation similar to (3) was first considered by Tang.[4]

[1]I. M. Begg and R. A. Cairns, J. Phys. D: Appl. Phys. **9**, 2341 (1976).
[2]C. J. Randall, J. J. Thomson, and K. G. Estabrook, Phys. Rev. Lett. **43**, 924 (1979).
[3]L. M. Gorbunov, Usp. Fiz. Nauk **109**, 631 (1973) [Sov. Phys. Usp. **16**, 217 (1973)].
[4]C. L. Tang, J. Appl. Phys. **37**, 2945 (1966).
[5]L. M. Gorbunov and A. N. Polyanichev, Zh. Eksp. Teor. Fiz. **74**, 552 (1978) [Sov. Phys. JETP **47**, 290 (1978)].
[6]H. E. Heuy, A. Mase, C. Pawley, T. Ohnuma, N. C. Luhmann, J. J. Thomson, and W. F. DiVergilio, Proc. Int. Conf. on Plasma Phys. Nagoya, Japan, 1980, Vol. 1, 7P-1-24.
[7]A. Mase, H. Heuy, M. Rhodes, J. Holt, and N. C. Luhmann, *ibid.*, 9-P-01.
[8]H. Akiyama, O. Matsumoto, and S. Takeda, *ibid.*, 10-P-11--08.
[9]A. Ng, L. Pitt, D. Saltzmann, and A. A. Offenberger, Phys. Rev. Lett. **42**, 307 (1979).
[10]I. L. Fabelinskiĭ, FIAN Preprint No. 143, 1972.
[11]A. V. Vinogradov, B. Ya. Zel'dovich, and I. I. Sobel'man, Pis'ma Zh. Eksp. Teor. Fiz. **17**, 271 (1973) [JETP Lett. **17**, 195 (1973)].

Translated by J. G. Adashko