EFFECT OF VERY SHORT LASER PULSES ON ABSORBING SUBSTANCES

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Evaporation of a solid body absorbing a very short duration laser pulse is considered. Evaporation occurs on relaxation of the strain resulting from compression by a strong shock wave moving in the solid body from the light-absorbing surface layer. Prior to the start of the hydrodynamic motion energy is transferred within the body by thermal waves. The law of motion of the thermal wave determines the mass of matter to which the absorbed energy is transferred and also the intensity of the produced shock wave. For not very small values of absorbed energy density Q, the mass of evaporated matter is proportional to $M \sim Q^{1-\alpha}$ with $\alpha < 1/6$.

A number of recent communications have reported the generation of very short (picosecond) high-power laser pulses (see, for example,[1]). The processes of interaction of such pulses with matter greatly differ from the well investigated processes[2-3] caused by laser pulses of ordinary duration. In particular, the mechanism of evaporation of the absorbing condensed substances and the dynamics of the spreading of the produced plasma are entirely different. We present below a qualitative analysis of these processes and estimate the mass of the evaporated matter and the effective per unit energy of evaporation. The estimate shows that at a specified absorbed energy Q, the evaporated mass is proportional to the quantity $M \sim Q^{1-\alpha}$, with $\alpha < 1/6$. Thus, the energy consumed in the evaporation of a unit mass at small values of τ increases very slowly with increasing energy density Q, like $\lambda \sim Q^{\alpha}$. For other irradiation regimes, the growth of the effective specific energy of evaporation with increasing Q is faster[2,3].

We shall consider radiation pulses whose duration τ is much shorter than the characteristic time τ_0 of the hydrodynamic motion. For such pulses it is possible to neglect the motion of the medium during the time of action of the radiation, and thus it is possible to consider separately the hydrodynamic and optical parts of the problem. The subsequent process proceeds in the following manner. The energy absorbed in the surface layer is transferred to the cold matter in a thermal wave propagating via nonlinear thermal conductivity. As a result, at the end of the pulse a certain mass M* will have a specific energy of the order of Q/M*. When this mass expands, the body acquires a momentum I on the order of $\sqrt{QM^*}$, and a strong shock wave propagates through the matter. The main evaporation occurs when the matter compressed by the shock wave relaxes. Thus, calculation of the evaporated mass reduces essentially to establishment of the law of motion of the shock wave.

The motion of the medium at a time $t > \tau$ is described by the well known "short impact" theory[4,5]. The limiting motion is self-similar and depends on the variable $\xi = x/X(t)$, where $X(t) = At^s$ is the coordinate of the shock-wave front (for concreteness we shall speak throughout of one-dimensional motion of the medium). In the case of "short impact" the exponent s is not determined, as is well known, by the conservation laws, but depends on the thermodynamic properties of the medium and should be determined by solving the equations of motion. For an ideal gas, a number of values of s as a function of the adiabatic exponent γ are listed in[5].

The intensity of the shock wave is determined by the value of the constant A. For its calculation, it is necessary, strictly speaking, to solve the non-self-similar problem describing the initial stage of the "short-impact." However, if we confine ourselves to an estimate of the quantity A, it suffices to estimate a combination of the momentum of the energy, in the form $Q^{2S-1}I^{2-3S}$; this combination is conserved in self-similar motion and is connected with the mass M* of the matter encompassed by the thermal wave at the start of the hydrodynamic motion.

In order to find the law of propagation of the thermal wave, it is necessary to specify the coefficient of thermal conductivity as a function of the temperature. We shall consider by way of a model of the absorbing body, at high temperatures, a fully ionized gas. It can be shown that for the values of the absorbed energy Q of practical interest (which are not too small), the decisive mechanism of energy transport in the plasma is the electronic specific heat. The energy transported by the radiation in cases of practical interest plays a secondary role, since the plasma layer containing a mass M* is thin in such cases. This means that the loss to radiation from the heated layer will also be small compared with the absorbed energy (the corresponding estimate is given below).

The coefficient of electronic thermal electronic activity is given by the formula[6]

$$\varkappa_e = BT^{5/2},\ B \approx 2\cdot 10^{-4}(Z \ln \Lambda)^{-1}\ \text{Oe/cm-sec-deg}^{7/2}.$$

For such $\kappa(T)$, the front of the thermal wave can be readily seen to propagate in accordance with the law

$$x \approx B\left(\frac{Q}{\theta}\right)^{5/9}\left(\frac{t}{Bc\rho}\right)^{7/9}, \quad \theta = \begin{cases} \tau, t < \tau \\ t, t \geqslant \tau \end{cases} \tag{1}$$

where $c\rho$ is the specific heat per unit volume. The average temperature behind the front of the wave is then

$$T \approx \left(\frac{Q}{\theta}\right)^{4/9}\left(\frac{t}{Bc\rho}\right)^{2/9}. \tag{2}$$

Calculating the mass of the matter heated by the thermal wave by the instant τ_0, we get

$$M^* \approx c^{-1}(Bc\rho\tau_0)^{2/9}Q^{5/9}. \tag{3}$$

Let us now estimate the radiation loss. For the bremsstrahlung mechanism, the energy radiated per unit volume per unit time is[5] $\epsilon - 1.4 \times 10^{-27} n_e^2\sqrt{T}$, and the total radiated energy is in any case smaller than $Q_1 = \epsilon t x(t)$.

Using formulas (1) and (2), we can write the condition $Q_1 \ll Q$ in the form

$$\tau \leqslant 4\cdot 10^{-13}(Bc\rho)^{4/5}Q^{1/5}, \quad Q \leqslant 10^{-15}c^{5/2}/B^2 \tag{4}$$

For the electronic density corresponding to a condensed body and for $Q = 10^{10}$ e/cm^2, we get from formula (4) $\tau \ll 3 \times 10^{-9}$ sec.

At small values of Q, even for very short pulses, an important role is assumed by radiation loss. When the loss is large, the law of motion of the thermal wave is radically altered: the sharp cooling of the surface layer gives rise to the thermal wave, described in[5], from the dipole source. This case, however, is of little practical interest and is not considered here.

In order to be able to regard a plasma layer as optically thin, it is necessary to satisfy the inequality $l_1 = x(\tau_0)$, where $l_1 = 1.5 \times 10^{23}T^{1/2}/n_e^2Z^2$ is a suitably[5] averaged radiation mean free path. From this inequality follows a limitation on the pulse energy:

$$Q \geqslant 10^{17}B^3 \sim 10^4 \text{ Oe/cm}^2.$$

It is easy to verify that the condition for applicability of the approximation of the electric thermal conductivity remains satisfied at the same time.

After the light pulse stops, the motion of the thermal wave slows down sharply. Further energy transport from the heated region is determined mainly by hydrodynamic motion, which becomes noticeable after a time lapse on the order of

$$\tau_0 = \frac{x(\tau_0)}{a} \approx \frac{Q}{\rho c^{3/2}}\left(\frac{Bc\rho\tau_0}{Q^2}\right)^{1/3},$$

where a is the speed of sound. Thus, the condition for applicability of the foregoing analysis, $\tau \ll \tau_0$, is given by the inequality

$$\tau \leqslant (BQ)^{1/2}/\rho c^{7/4}. \tag{5}$$

In the already considered case of the electron density of a solid and $Q = 10^{10}$ Oe/cm^2, we get from (5) $\tau \ll 10^{-9}$ sec.

In the case of very short pulses, an important role may be played by the violation of equilibrium between the electrons and the ions. It is obvious that the relaxation time τ_{ei} is the lower limit for τ_0; for pulses whose duration is $\tau < \tau_{ei}$, we get the above-described qualitative picture, the only difference being that the thermal wave propagates over the electrons with the ions cold. This case differs from the equilibrium case only in the values of the constants, but this cannot change noticeably the obtained estimates.

To determine the mass of matter evaporated in the rarefaction wave, we note that in accordance with[5] complete evaporation takes place under the condition that the energy density in the body compressed by the shock wave exceeds approximately by a factor of 5 the heat of the phase transition: $E > 5\mathscr{L}$. Calculating the constant A by the method described above, using formula (3), we can readily obtain for the evaporated mass the expression

$$M = \frac{Q}{\lambda}, \quad \lambda = \lambda_0\left(\frac{Q\sqrt{\lambda_0}}{T_0\varkappa(T_0)}\right)^\alpha, \tag{6}$$

where $\alpha = (2 - 3s)/6(1 - s)$ and $\lambda_0 = RT_0 = 5\mathscr{L}$. Since for s we have the inequality[5] $1/2 < s < 2/3$, the exponent α turns out to be, in any case, smaller than 1/6.

In estimating the quantity λ we did not take into account the loss to radiation during the stage of the shock wave, i.e., the motion of the matter in the "short impact" is assumed to be adiabatic. To estimate the energy loss, we note that, for values of Q and τ of practical interest, the layer of matter encompassed by the shock wave becomes opaque very rapidly, owing to the sharp temperature dependence of the radiation mean free path. It is easy to see, further, that the main contribution to the energy loss is made by emission from the layer prior to the loss of the transparency. The corresponding energy is certainly much smaller than the quantity

$$Q_2 = \varepsilon(T_0)X_0t_0(l_{10}/X_0)^{2s/(7-6s)},$$

where the subscript zero designates quantities corresponding to the energy density behind the shock wave, equal to $5\mathscr{L}$. From the inequality $Q_2 \ll Q$, as can be readily seen, follows also one more limitation on the energy of the laser pulse. Thus, in the case already considered, Q should be much smaller than 10^{11} Oe/cm^2.

It follows from (6) that in the regime under consideration the effective value of the specific heat of evaporation at small values of τ increases very slowly with increasing Q. We note that, for example, in the regime of self-similar expansion of the absorbing plasma with constant optical thickness[2], the increase of τ with increasing Q turns out to be faster: $\lambda \sim Q^{1/2}$. The reason for the difference can be readily understood if account is taken of the fact that under conditions of a "short impact" the main mass of the evaporated matter has a specific energy on the order of λ_0 even at very large values of Q.

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