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VAPORIZATION OF MATTER EXPOSED TO LASER EMISSION

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The process of vaporization of matter exposed to laser emission is considered theoretically within light flux density range of 10^6—10^9 W/cm^2. The analysis is based on gas-dynamics equations. It is shown that in the case of transparent material two gas-dynamic modes of vaporization and heating of the material are possible, depending upon the radiation flux density. The first mode corresponds to low flux densities when the temperature of the radiation-absorbing layer of condensed matter is below critical temperature and a phase transformation of condensed matter into gas takes place. The other mode relates to fluxes capable of transforming condensed matter into gas by expansion due to thermal pressure. Boundary conditions are formulated and exact analytical solutions of the gas-dynamic equations are derived to describe the process of vaporization and heating of matter in both cases. Expressions for the gas-dynamic quantities are derived as functions of the radiation flux density q^0 and the parameters of condensed matter.

THE investigation of processes accompanying the action of powerful laser emission on matter has recently been the subject of considerable attention. A large number of papers (see review paper[1]) deal with the phenomenon of gas breakdown caused by focused laser emission. There is also a number of papers, mainly of experimental nature, investigating the effect of powerful laser emission on solids (see review paper[2]). Nevertheless, there are only a few theoretical papers investigating the processes of interaction between laser emission and solids. Ready[3,4] presented a qualitative analysis of two possible mechanisms for the interaction of laser emission with solids, depending on the magnitude of the incident flux density. When the flux densities are small, the layer of radiation-absorbing material undergoes a phase transformation at a temperature which the author calls the "temperature of vaporization." When the flux density reaches $\sim 10^8$ W/cm^2, the temperature of the material rises above critical and there is no phase transformation. A quantitative analysis of these processes presented in the above papers is based on the thermal conductivity equation for the condensed phase in a manner similar to[5].

Anisimov, Bonch-Bruevich, and others[5] relate the thermal conductivity mechanism to the boundary condition of free vaporization of the solid into vacuum.[6] The use of the activation formula[6] seems to yield a good approximation of the vaporization process for fluxes at which the density of the vaporized matter is low and its temperature is considerably below the critical point.

It is obvious, however, that thermal conductivity plays a significant role only when the flux density is very low and the radiation absorbed during a pulse goes to heat rather than vaporize the solid. Otherwise, when the layer vaporized during a pulse is thicker than that heated by thermal conductivity, the thermal conductivity mechanism is significant only during the initial stage and the total energy balance may include only the energy expended on vaporization. The thermal conductivity mechanism is then responsible only for the distribution of temperature in the condensed phase beyond the vaporization limit.

The minimum flux density q' that can cause the vaporization process to play the principal role can be obtained from the condition that the internal energy of the layer spanned by the thermal wave becomes comparable with the specific heat of vaporization Q during the laser emission pulse τ. This yields the expression $q' \approx \Omega \rho_0 \sigma^{1/2} \tau^{1/2}$ where ρ_0 is density of the condensed matter and σ is the coefficient of temperature conductivity.

We consider here the range of radiation fluxes $q^0 > q'$ at which, as noted above, the internal energy of the radiation-absorbing layer exceeds the binding energy Ω, thus leading to vaporization, i.e., to gas-dynamic motion of matter.

A similar situation occurs in magneto-implosive generators[18] and in gas-discharge light sources.

A number of gas dynamic problems can also be found in the book by Zel'dovich and Raĭzer[7]. The analytic solution method developed below is analogous to that suggested by Sakharov[18].

We consider a one-dimensional plane problem in which the radiation is incident at the time t = 0 along the X axis from the x = ∞ side on a solid surface normal to and intersecting the X axis at the point x = 0. Using the notation of[8], we have

$$\frac{\partial \rho}{\partial t} + \frac{\partial}{\partial x}(\rho v) = 0,$$
$$\frac{\partial}{\partial t}(\rho v) + \frac{\partial}{\partial x}(p + \rho v^2) = 0,$$
$$\frac{\partial}{\partial t}\left(\rho\varepsilon + \rho\frac{v^2}{2}\right) + \frac{\partial}{\partial x}\left[\rho v\left(\varepsilon + \frac{v^2}{2} + \frac{p}{\rho}\right)\right] + \frac{\partial q}{\partial x} = 0, \quad (1)$$

where v, ρ, p, and ϵ are the velocity, density, pressure, and internal energy of the vapor respectively.

In the energy equation, the term $\partial q/\partial x$ defines the absorption of radiation both in a thin layer l on the surface of the condensed matter ($l \sim 1/K_0$, where K_0 is the absorption coefficient, $\sim 10^4$—10^5 cm^{-1} for strongly absorbing materials such as metals) and in the vaporized layer. However, the occurrence of absorption by the vapor depends on the incident flux q^0. In fact, a sharp transition from condensed to gaseous state takes place in the thin absorbing layer l at the solid surface. Such a transition is accompanied by a steep decrease of the absorption coefficient if the vapor temperature is low in comparison with the ionization potential. In this case we have a vaporization mode in which the vapor is transparent, i.e., the $\partial q/\partial x$ term of the energy equation (1) can be dropped everywhere except for a thin layer at the surface of the condensed matter. The character of the vaporization in this case significantly depends upon the processes occurring within the absorbing layer. When the radiation flux is small we must take phase transformation on the condensed surface into account and the phase interface is the boundary of the absorbing layer so long as vapor density is sufficiently low. As the flux increases, however, the vapor density at the phase interface can become large enough to render the vapor opaque to the incident radiation. In such a case the absorption boundary shifts into the gaseous phase and its location is determined by the particular value of density that renders the material opaque. Consequently, if this density ρ^* is less than critical density $\rho_{cr} \approx \rho_0/3$, then the phase transformation must be accounted for only until the vapor density at the condensed surface reaches the value of ρ^*. On the other hand, if $\rho^* \approx \rho_{cr}$, the phase transformation will occur at the boundary of the absorbing at incident fluxes up to those for which the temperature of the condensed body is $T_0 \approx T_{cr}$. When $\rho^* > \rho_{cr}$ the absorption boundary coincides with the phase interface for flux values corresponding to a surface temperature T_0 lower than the boiling point at the density ρ^*. Further increase in flux shifts the absorption boundary into the condensed phase, whose transparent layer expands adiabatically in a manner similar to the relaxation of material subjected to a shock wave emerging on a solid surface.[7]

Increasing the incident flux q^0 causes the vapor temperature to increase and its absorption coefficient is increased by ionization and excitation of the atoms. This, in turn, leads to a new mode of vaporization in which the absorption of radiation by the vapor is fully responsible for the dynamics of the entire process. In this case the term $\partial q/\partial x$ in (1) is different from zero over the entire region subject to the motion. For the absorption coefficient of a plasma this problem was considered by Nemchikov, Krol' and the present authors[9], and also by Caruso, Bertotti, and Giupponi[10].

A characteristic feature of the vaporization of material by laser emission is the presence of a thin layer, of $\sim 10^{-4}$—10^{-5} cm, on the surface of the condensed body; this layer is a region of strong absorption, $\partial q/\partial x \sim q_0/l$, and consequently a region of sharp variation of all gas-dynamic quantities. Regarding it as a region of marked discontinuity, we write the conditions of conservation of mass flow, momentum, and energy within the discontinuity

$$\rho_1(D - v_1) = \rho_0 D, \quad p_1 - \rho_0 D v_1 = p_0,$$
$$-\rho_0 D(\varepsilon_1 + v_1^2/2) + p_1 v_1 = q_0, \quad (2)$$

where D is the velocity of vaporization boundary (D < 0), v_1, ρ_1, and p_1 are respectively the velocity, density, and pressure of the gas at the interface with the condensed body, p_0 and ρ_0 are the pressure and density respectively in the condensed body, q_0 is the flux incident on the surface of the condensed body, $\epsilon_1 = p_1/(\kappa - 1)\rho_1 + \Omega$, and κ is the adiabatic index. Equations (2) are analogous to the relations on the front of a detonation wave[11,12] and were derived by Raĭzer[13] for the case of gas breakdown induced by focused laser emission. They can be considered as boundary conditions of the system defined by (1).

1. We first consider the case when the incident radiation flux q^0 is in the range $q' < q^0 < q''$, where q'' is the lower limit of flux values for which there is no phase transformation at the absorption boundary. As noted above, the vapor is transparent in this region and the phase transformation at the surface of the condensed body must be taken into account.

A rigorous solution of this problem calls for an analysis of the kinetics of phase transformation. The condensation rate in supersaturated vapor was considered by Volmer[14], Frenkel[15], Zel'dovich[16], and Raĭzer[17]. The mass increase of the condensed phase was determined in these papers by the flow of vapor particles onto the surfaces of nuclei whose number and critical dimensions depended on the degree of supersaturation.

In our case there could take place in the absorbing layer a "volume" mechanism of vapor formation, due to the flow of particles from the superheated liquid to vapor nuclei whose dimensions and quantity are determined by the degree of superheat. The distinguishing feature of our problem, however, is that the interface between the two phases can also be a site of vapor formation. The rate of the "volume" vaporization becomes comparable to the rate of vaporization at the phase interface when the total surface area of the nuclei of critical size, which are centers of stable stationary vaporization, is equal to the area of the interface. According to[6,15], the number of nuclei of critical size r per unit volume is

$$N \approx n_0 \exp\{-4\pi\alpha r^2/3kT\}, \tag{3}$$

where n_0 is the particle-number density in the condensed phase and α is the coefficient of surface tension. Then the ratio of the total area Σ of the nuclei in the absorbing layer l to the free surface area S is given by the expression

$$\frac{\Sigma}{S} = 4\pi r^2 n_0 l \exp\left\{-\frac{4\pi\alpha r^2}{3kT}\right\}. \tag{4}$$

The maximum of Σ/S occurs when

$$r^2 = r_1^2 = 3kT/4\pi\alpha, \tag{5}$$

and $\Sigma/S = 1$ when $r^2 = r_2^2 > r_1^2$.

Vapor pressure p_v in the nucleus is related to the pressure p_0 in the condensed medium by the equation

$$p_v = p_0 + 2\alpha/r. \tag{6}$$

The temperature T_0 corresponding to phase equilibrium across a plane interface at the pressure p_0 is determined by the equation of the phase-equilibrium curve

$$p_0 \approx \frac{BR_0T_0}{\mu}\exp\left\{-\frac{\omega}{kT_0}\right\}, \tag{7}$$

where B ≈ const, R_0 is gas constant, μ is the gram-atomic weight, and ω is the heat of vaporization per atom. The degree of superheat of the medium is defined by the ratio of the condensed-phase temperature T to the temperature T_0. Combining (6) with the equation of the phase equilibrium curve with allowance for the curvature[6,15]

$$p_v = \frac{BR_0T}{\mu}\exp\left\{-\frac{\omega}{kT} - \frac{2\alpha}{n_0 rkT}\right\}, \tag{8}$$

we obtain an equation for the degree of superheat under the condition $\Sigma/S = 1$:

$$\frac{BR_0T}{\mu}\exp\left\{-\frac{\omega}{kT} - \frac{2\alpha}{n_0 r_2 kT}\right\} = \frac{BR_0T_0}{\mu}\exp\left\{-\frac{\omega}{kT_0}\right\} + \frac{2\alpha}{r_2}. \tag{9}$$

In the actual case considered here, with a thin ($l \sim 10^{-4}$—10^{-5} cm) absorbing layer in which the volume vapor formation process can take place, the terms associated with the nuclear-surface curvature play the principal role, according to a numerical analysis of (9). As a result, the condition $\Sigma/S \approx 1$ is reached at temperatures $T \sim 0.3\,\omega/k$, i.e., when the temperature of the condensed phase is close to critical (the following numerical values were used in the analysis: $l = 10^{-4}$ cm, $n_0 = 10^{23}$ cm^{-3}, $\alpha = \omega n_0^{2/3}$ [15], $r_2^2 \approx 7r_1^2 \approx 2kT/\omega n_0^{2/3}$, and B selected so as to have $B \approx 0.3\rho_0 e^{\omega/kT}$ when $T \approx T_{cr}$).

Consequently, the degree of superheat T/T_0 necessary for the transition into the volume vaporization mode when $T_0 \ll \omega/k$ turns out to be high in the case under consideration; at the same time, when $T_0 \ll \omega/k$, i.e., when phase transformation must be taken into account, vaporization at the phase interface plays the principal role.

The vaporization mechanism substantially depends on the characteristic time of establishment of the phase equilibrium. If this time is insufficient, the phase transformation will be accompanied by "evaporation" with bond rupture due to the expansion of the material under thermal pressure.

We assume that phase equilibrium is established in a time period that is short enough in comparison with l/D, so that a phase transformation occurs in the incident-flux absorption region at the saturated vapor pressure $p_0 = p_1 - \rho_0 D v_1$. This means that effective vapor formation occurs at the boiling point at a pressure corrected for the recoil effect; the contribution of the recoil effect to the pressure is $-\rho_0 D v_1$.

The system of equations (1) with $\partial q/\partial x = 0$ and boundary conditions (2) and (3) has when $q^0 = q_0 = \text{const}$ a solution that is the limiting case of a centered rarefaction wave[7,8] corresponding to an isentropic expansion of gas in vacuum. The centered rarefaction wave problem is a self-similar problem, so that all the gas-dynamic quantities are functions of a single variable $\xi = x/t$. A dimensionless variable $\lambda = \Omega^{1/2}x/t$ is convenient in this case. We note that the selection of such a variable is not unique, since only any two of the three key parameters q_0, Ω, and ρ_0 have independent dimensions in our case. This circumstance, however, will not affect the end result. The unknown functions v, ρ, and p can be represented in the form

$$v = \Omega^{1/2}V(\lambda), \quad \rho = q_0\Omega^{-3/2}R(\lambda), \quad p = q_0\Omega^{-1/2}P(\lambda), \tag{10}$$

where V(λ), R(λ), and P(λ) are the dimensionless velocity, density, and pressure respectively.

The system of equations (1) will then have the form

$$\frac{d}{d\lambda}(RV) - \lambda\frac{dR}{d\lambda} = 0, \quad (V-\lambda)\frac{dV}{d\lambda} + \frac{1}{R}\frac{dP}{d\lambda} = 0,$$
$$(V-\lambda)\frac{d}{d\lambda}\left(\frac{P}{R}\right) + (\varkappa-1)\frac{P}{R}\frac{dV}{d\lambda} = 0. \tag{11}$$

The solution of (11) can be represented as follows:

$$R = R(\lambda_1)\left[1 - \frac{\lambda-\lambda_1}{\lambda_2-\lambda_1}\right]^{2/(\varkappa-1)},$$
$$P = R(\lambda_1)\frac{(\varkappa-1)^2(\lambda_2-\lambda_1)^2}{\varkappa(\varkappa+1)^2}\left[1 - \frac{\lambda-\lambda_1}{\lambda_2-\lambda_1}\right]^{2\varkappa/(\varkappa-1)}, \tag{12}$$
$$V = \left[\frac{(\varkappa-1)\lambda_2}{\varkappa+1} + \frac{2\lambda_1}{\varkappa+1}\right]\left[1 + \frac{\lambda-\lambda_1}{(\varkappa-1)\lambda_2/2-\lambda_1}\right],$$

where λ_2 and λ_1 are the values of the self-similar variable λ corresponding respectively to the interfaces of the vaporized material with the vacuum and with the surface of the condensed body ($\lambda_2 > 0$ and $\lambda_1 < 0$).

The solutions (12) yield at the vacuum interface

$$P(\lambda_2) = 0, \quad R(\lambda_2) = T(\lambda_2) = 0,$$

where T is the gas temperature. In addition, the solutions (12) satisfy the Jouguet condition at the condensed-body interface:

$$v_1 - c_1 = D, \tag{13}$$

where $C_1 = (\kappa p_1/\rho_1)^{1/2}$ is sound velocity in the condensed body gas at the interface. Condition (13) means that the values of gas-dynamic quantities at the moving interface between the condensed body and the gas are stationary and we have the following boundary characteristic for the centered rarefaction wave in our case:

$$x_0 = Dt. \tag{14}$$

Equation (13) together with the boundary conditions (2) and the equation of the phase-equilibrium curve comprise a system of five equations with six unknown parameters p_0, p_1, v_1, D, ρ_1, and ρ_s, where ρ_s is the saturated-vapor density. The sixth equation needed to close the system must contain the relation between the temperatures of vapor and the condensed phase. It is simplest to assume equal phase temperatures at the interface. This yields the condition

$$p_0/\rho_s = p_1/\rho_1. \tag{15}$$

Following[15], we write the equation of the phase curve

$$\rho_s = B\exp\{-\Omega\rho_s/p_0\}, \tag{16}$$

where B can be considered constant.

Substituting (10) and (12) into the first and third equations of (2) and combining the second equation of (2) with (15) and with the equation of the phase equilibrium curve, we obtain a system of three equations with respect to three unknowns λ_1, λ_2, and $R(\lambda_1)$:

$$\eta R(\lambda_1) = (\varkappa+1)\lambda_1/(\varkappa-1)(\lambda_1-\lambda_2),$$
$$\frac{(\varkappa+1)^2\lambda_1[\lambda_1+(\varkappa-1)\lambda_2]}{(\varkappa-1)^2(\lambda_1-\lambda_2)^2} = -\beta\exp\left\{-\frac{\varkappa(\varkappa+1)^2}{(\varkappa-1)^2(\lambda_2-\lambda_1)^2}\right\},$$
$$\lambda_1\left[1 + \frac{\lambda_1^2}{\varkappa(\varkappa+1)} + \frac{(\varkappa-1)\lambda_2^2}{2(\varkappa+1)} + \frac{\lambda_1\lambda_2(\varkappa-1)}{\varkappa(\varkappa+1)}\right] = -\eta, \tag{17}$$

where $\eta = q_0/\rho_0\Omega^{3/2}$ and $\beta = B/\rho_0$.

The solution of the system of algebraic equations (17) completely defines the given problem. In principle this system can be solved numerically; however, the obvious condition $|\lambda_1| \ll \lambda_2$ significantly simplifies the equations and yields expressions for the gas dynamic quantities in finite form. The above condition renders the last two equations of (17) equivalent to

$$\frac{(\varkappa-1)^2\lambda_2^3}{2(\varkappa+1)^3} + \frac{(\varkappa-1)\lambda_2}{(\varkappa+1)^2} = \frac{\eta}{\beta}\exp\left\{\frac{\varkappa(\varkappa+1)^2}{(\varkappa-1)^2\lambda_2^2}\right\},$$
$$\lambda_1 = -\frac{\beta\lambda_2(\varkappa-1)}{(\varkappa+1)^2}\exp\left\{-\frac{\varkappa(\varkappa+1)^2}{(\varkappa-1)^2\lambda_2^2}\right\}. \tag{18}$$

The system (18) determines the unknown quantities λ_1 and λ_2 as functions of the parameters $\eta = q_0/\rho_0\Omega^{3/2}$ and $\beta = B/\rho_0$. An analysis of (18) shows that the expressions for λ_1 and λ_2 can be represented within a wide range of the parameter η (2–3 orders of magnitude) in the form:

$$\lambda_2 = \left(\frac{\gamma}{A - \ln(\eta/\beta)}\right)^{1/2}, \quad \lambda_1 = \frac{C(\varkappa-1)\eta}{(\varkappa+1)^2}, \tag{19}$$

where $\gamma = \kappa(\kappa+1)^2/(\kappa-1)^2$ and A and C vary little with the parameter η; their approximate values can be determined, depending upon the region of variation of this parameter.

We now write explicit approximate expressions for the gas-dynamic quantities at the interface with the condensed body.

Taking into account (10) and the solutions (12) and (19), and assuming $\lambda = \lambda_1$, we obtain

$$\rho_1 = \frac{(\varkappa+1)\rho_0|\lambda_1|}{(\varkappa-1)\lambda_2} \approx Cq_0/(\varkappa+1)\Omega^{3/2}\left(\frac{\gamma}{A-\ln(\eta/\beta)}\right)^{1/2},$$

$$v_1 = \frac{(\varkappa-1)\Omega^{1/2}\lambda_2}{\varkappa+1} \approx \frac{(\varkappa-1)\Omega^{1/2}}{\varkappa+1}\left(\frac{\gamma}{A-\ln(\eta/\beta)}\right)^{1/2},$$

$$D = \Omega^{1/2}\lambda_1 \approx -C(\varkappa-1)q_0/(\varkappa+1)^2\rho_0\Omega,$$

$$p_1 = \frac{(\varkappa-1)|\lambda_1|\lambda_2\rho_0\Omega}{\varkappa(\varkappa+1)} \approx \frac{C(\varkappa-1)^2 q_0}{\varkappa(\varkappa+1)^3\Omega^{1/2}}\left(\frac{\gamma}{A-\ln(\eta/\beta)}\right)^{1/2}$$

$$T_1 = \frac{(\varkappa-1)^2\mu\Omega\lambda_2{}^2}{\varkappa(\varkappa+1)^2 R_0} \approx \frac{\mu\Omega}{R_0(A-\ln(\eta/\beta))} \qquad (20)$$

As an example we give numerical results obtained for iron ($B = 3.3 \times 10^4$ g/cm^3[17], $\Omega = 6.9 \times 10^{10}$ erg/g). Given $\kappa = 5/3$ and a range of flux densities $q_0 \approx 10^6$–10^9 W/cm^2, the values of A, C, and γ are in this case 1.9, 8, and 26.7 respectively.

Using the solutions (20), we can determine the flux density q''. Assuming that the vapor density at which the vapor becomes opaque is $\rho^* \lesssim \rho_{cr}$, we find from the first equation of (20)

$$q'' \simeq \frac{(\varkappa+1)\gamma^{1/2}\rho^*\Omega^{3/2}}{C(A-\ln(\eta/\beta))^{1/2}}, \qquad (21)$$

where $(A - \ln(\eta/\beta))$ can be defined at the point where $T_1 \approx T_{cr}$ by making use of the last equation of (20).

2. We now consider a case in which the incident flux q^0 is limited to the range $q'' < q^0 < q'''$; q''' is the flux at which the thermal energy of the atoms of the vaporized material approaches the ionization potential. Consequently, any further increase in the flux $q^0 > q'''$ makes it necessary to take radiation absorption by the vapor into account.

In this case the material in the thin absorbing layer is a strongly interacting gas of high density. Thermal pressure causes the material to expand and to become transparent to incident radiation at the density ρ^*. Further gas-dynamic motion of the transparent material will clearly be the same as in the above case, i.e., it will be defined by a centered rarefaction wave.

The introduction of a specified density $\rho^* = \rho_1$ into the conditions of the problem provides a substantial simplification. In this case the three equations of (2) and Jouguet condition (13) completely determine four unknown quantities: p_0, p_1, v_1, and D.

The density ρ^* can in principle be determined if we know the coefficient of absorption as a function of ρ and T. However the dependence of the absorption coefficient on ρ and T in the range of densities comparable to the initial density ρ_0 is not known. Nevertheless, in the gas-dynamic problem under consideration the quantity ρ^* is an external parameter that does not affect the nature of the gas-dynamic motion. In this connection, the concrete value necessary to compute the gas-dynamic quantities can be obtained from experiments[19]. The data on exploding wires[20] can be used for metals where the conductivity (and consequently the absorption coefficient) turns to zero at a definite density value. A theoretical evaluation can also be carried out by considering the transition of metal into dielectric (the Mott transition[21,22]). In this case, the density ρ^* is determined by the condition that the Debye length equal the Bohr radius a_0, i.e., when the electron density satisfies the relationship $n_e^{-1/3} = 4a_0$.

The first and third equations of (2), taking into account (10) and the solutions (12), assume in this case the following form:

$$\lambda_2 = \lambda_1[1-(\varkappa+1)/\delta(\varkappa-1)], \qquad F\lambda_1{}^3 + \lambda_1 + \eta = 0, \quad (22)$$

where $\delta = \rho^*/\rho_0$ and

$$F = \frac{1}{\delta^2}\left[\frac{2-\varkappa}{\varkappa(\varkappa-1)} + \frac{1}{2}\left(\delta^2 - \frac{2\delta(\varkappa-1)}{\varkappa} + 1\right)\right].$$

Furthermore, it follows from (6) that

$$R(\lambda_1) = \delta/\eta \qquad (23)$$

We write explicit expressions for the gas-dynamic quantities at the interface with the condensed body. Considering (10), (12), (22), and the first equation of (22), we obtain

$$\rho_1 = \delta\rho_0, \quad p_1 = \frac{\Omega\rho_0\lambda_1{}^2}{\varkappa\delta}, \quad v_1 = \Omega^{1/2}\frac{(\delta-1)\lambda_1}{\delta},$$

$$D = \Omega^{1/2}\lambda_1, \quad T = \frac{\mu\Omega\lambda_1{}^2}{R_0\varkappa\delta^2}. \qquad (24)$$

The self-similar constant λ_1 is determined by the second equation of (22), which has only one real root equal to

$$\lambda_1 = \left\{\left[\left(\frac{\eta}{2F}\right)^2 + \frac{1}{27F^3}\right]^{1/2} - \frac{\eta}{2F}\right\}^{1/2}$$

$$-\left\{\left[\left(\frac{\eta i}{2F}\right)^2 + \frac{1}{27F^3}\right]^{1/2} + \frac{\eta}{2F}\right\}^{1/3}. \qquad (25)$$

Our problem is completely solved by the substitution of (25) into (24).

The exact values of the gas-dynamic quantities (24) and (25) depend in a complex manner on the parameter η, i.e., on the density of the incident flux q^0. It is therefore of interest to consider the limiting cases. A comparison with (21) readily shows that the flux $q^0 = q''$ corresponds to the parameter value $\theta \equiv \eta F^{1/2} = 1$. Therefore expressions (24) and (25) are valid within the parameter range $\theta > 1$. Nevertheless, if the establishment of phase equilibrium takes a long time, as it was noted above, (24) and (25) hold even when $\theta < 1$.

In this case, retaining the squared terms in the expansion of (25) in terms of θ, we obtain

$$\lambda_1 = -\eta\left(1-\frac{\eta F^{1/2}}{\sqrt{3}}\right), \quad v_1 = \frac{(1-\delta)q_0}{\delta\Omega\rho_0}\left(1-\frac{q_0F^{1/2}}{\sqrt{3}\Omega^{3/2}\rho_0}\right),$$

$$p_1 = \frac{1}{\varkappa\delta\rho_0}\left(\frac{q_0}{\Omega}\right)^2, \quad T_1 = \frac{\mu}{R_0\varkappa\delta^2}\left(\frac{q_0}{\rho_0\Omega}\right)^2,$$

$$D = -\frac{q_0}{\Omega\rho_0}\left(1-\frac{q_0F^{1/2}}{\sqrt{3}\rho_0\Omega^{3/2}}\right). \tag{26}$$

For a certain value of the incident flux density, within the range $\theta > 1$,

$$q^0 = 5\rho_0\Omega^{3/2}/F^{1/2}, \tag{27}$$

the specific thermal energy of the vaporized material becomes equal to the heat of vaporization Ω. We note that since the heat of sublimation per atom is usually comparable with the ionization energy, one cannot in principle neglect the absorption of radiation by the vaporized material at flux densities $q^0 > q''' = 5\rho_0\Omega^{3/2}/F^{1/2}$. As noted above, the self-consistent mode of heating and vaporization considered in[9,10] may occur in such a case.

It is possible, however, that the heat of sublimation is much lower than the energy of ionization in the vaporized material. Vapor absorption can then be neglected as before, and the problem under consideration is reduced to a regime that is independent of the specific heat of evaporation. In fact, if $\eta F^{1/2} \gg 1$, it follows from (25) that

$$\lambda_1 = -(\eta/F)^{1/2}. \tag{28}$$

The corresponding expressions for the gas-dynamic quantities at the solid interface have the form:

$$v_1 = \frac{(1-\delta)}{\delta F^{1/3}}\left(\frac{q_0}{\rho_0}\right)^{1/3}, \quad p_1 = \frac{\rho_0{}^{1/3}q_0{}^{2/3}}{\varkappa\delta F^{2/3}},$$

$$T_1 = \frac{\mu}{R_0\varkappa\delta^2F^{2/3}}\left(\frac{q_0}{\rho_0}\right)^{2/3}, \quad D = -\frac{1}{F^{1/3}}\left(\frac{q_0}{\rho_0}\right)^{1/3}. \tag{29}$$

The bulk of the radiation energy goes in this case to heat the material.

In conclusion we note that our problem covers a broad range of laser emission fluxes wherein the gas-dynamic motion of the vaporized material plays a significant role. The analytic expressions obtained for the gas-dynamic quantities can be used for comparison with experimental data. The formal solution based on (2), together with the assumption of isentropic gas flow in the form of a centered rarefaction wave, is analogous to the problem of detonation in a free-surface material[11,12]. In our case, however, the ratio of the total liberated energy to the mass in motion is a nonlinear function of the flux. In physical terms this means that we have a "division" of the total incident energy into parts corresponding to an increase of the mass of the vaporized material and to heating of the material, and that the division significantly depends upon flux q^0.

We also note that the assumption of equal phase temperatures, which we have adopted when analyzing the conditions of phase transformation, is not sufficiently justified. The exact relation between the temperatures of the vapor and the condensed phase can be established by analyzing the kinetics of vapor formation at the plane interface.

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[1] Yu. P. Raĭzer, UFN 87, 29 (1965), Soviet Phys. Uspekhi 8, 650 (1966).

[2] S. Panzer, Z. Ang. mat. und Phys. 16, 138 (1965).

[3] J. F. Ready, Proc. Electronics Cong. Chicago, 1964.

[4] J. F. Ready, J. Appl. Phys. 36, 462 (1965).

[5] S. I. Anisimov, A. M. Bonch-Bruevich, M. A. El'yashevich, Ya. A. Imas, G. S. Romanov, and N. A. Pavlenko, ZhTF 34, 7 (1966) [sic!]

[6] L. D. Landau and E. M. Lifshitz, Statisticheskaya fizika (Statistical Physics), Fizmatgiz, 1964.

[7] Ya. B. Zel'dovich and Yu. P. Raĭzer, Fizika udarnykh voln i vysokotemperaturnykh gidrodinamicheskikh yavleniĭ (Physics of Shock Waves and High Temperature Hydrodynamic Phenomena), Fizmatgiz, 1954, Chapts. 1 and 10.

[8] L. D. Landau and E. M. Lifshitz, Mekhanika sploshnykh sred (Mechanics of Continuous Media), Gostekhizdat, 1954, Secs. 82, 144, and 150.

[9] Yu. V. Afanas'ev, V. M. Krol', O. N. Krokhin, and I. V. Nemchinov, Prikladnaya matematika i mekhanika 7 (1966).

[10] A. Caruso, B. Bertotti, and P. Giupponi, Raporto Interno L. G. I., 65/29 bis.

[11] Ya. B. Zel'dovich and A. S. Kompaneets, Teoriya detonatsiĭ (Theory of Detonation), Gostekhizdat, 1955, Chapts. 1 and 4.

[12] K. P. Stanyukovich, Neustanovivshiesya dvizheniya sploshnoĭ sredy (Unsteady Motions of a Continuous Medium), Gostekhizdat, 1955, Secs. 42—44, and 50.

[13] Yu. P. Raĭzer, JETP 48, 1508 (1965), Soviet Phys. JETP 21, 1009 (1065).

[14] O. Volmer, Z. Elektrochemie 35, 555 (1929).

[15] Ya. I. Frenkel', Statisticheskaya fizika (Statistical Physics), AN SSSR, 1948.

[16] Ya. B. Zel'dovich, JETP 12, 525 (1942).

[17] Yu. P. Raĭzer, JETP 37, 1741 (1959), Soviet Phys. JETP 10, 1229 (1960).

[18] A. D. Sakharov, UFN 88, 725 (1966), Soviet Phys. Uspekhi.

[19] V. M. Kul'gavchuk, PTÉ **1**, 132 (1965) and **2**, 159 (1966).

[20] F. Webb, H. H. Hilton, P. H. Levine, and A. V. Tollestrun, Exploding Wires, v. 2, IIL, 1963, p. 47. [Plenum Press, N. Y. 1962].

[21] R. G. Arkhipov, JETP **43**, 349 (1962), Soviet Phys. JETP **16**, 251 (1963).

[22] N. Mott, Nuovo Cimento Suppl. **7**, 318 (1958).

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