QUASISTATIONARY SPHERICALLY SYMMETRIC FLOW OF AN INTENSIVELY EMITTING PLASMA HEATED BY LASER RADIATION

A. V. Dobkin, T. B. Malyavina,
and I. V. Nemchinov

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Estimates and numerical computations show that under a sufficiently prolonged action of laser radiation on a target in a vacuum under plane geometry conditions of vapor transit, a significant part of the laser pulse energy is converted into thermal radiation emitted by the plasma. If the vapor layer thickness during the laser action becomes commensurate with the characteristic size of the spot being irradiated, r_b, then the vapor density is reduced more rapidly in the lateral direction because of the jet broadening being initiated than in the case of transit in the form of a plane layer. The peripheral layers of the plasma torch become transparent and transmit the laser radiation practically totally, absorption is realized only near the target, and a quasistationary self-consistent heating and vapor transit mode is built up [5, 6]. The plasma parameters attained at the end of the plane transit stage later seem "to be frozen." Consequently, modes in which the thermal radiation becomes a dominant factor can be realized even in the stationary stage. Since the role of the thermal radiation grows as the pulse duration increases in the plane stage, as does the vapor layer thickness x, correspondingly, [1-4], then the radiation yield in the stationary stage grows as the spot size r_b increases [6].

The quasistationary mode of intensively emitting plasma motion, in which the energy of an external source is liberated, is of considerable interest since its mass flow rate, pressure, density, temperature, and optical plasma thickness as well as emitted radiation flux density, are sustained at a constant level. This facilitates investigation of the optical properties of an emitting plasma and the practical utilization of the emitted radiation. In the quasistationary mode the plasma mass, pressure pulse, and emitted radiation pulse are proportional to the energy supplied, i.e., they can, in principle, grow to arbitrarily large values.

Description of the radiation-gas dynamic processes occurring in a two-dimensional nonstationary or quasistationary motion of a radiating vapor jet is a rather complex problem. We limit ourselves here just to the simpler case of radially symmetric steady motion of radiating vapors. We make certain additional assumptions about the nature of the radiation.

1. We will describe the radiation transport in the quasivolume de-excitation approximation [7, 8]. Let us compute the radiation intensity I_ε being propagated in a uniformly heated volume of substance under different values of its temperature T, density ρ, and characteristic dimension L. We find the integrated spectral radiation intensity F_r on the volume boundaries

$$I_\varepsilon = B_\varepsilon(1-\exp(-\varkappa_\varepsilon m)),\ B_\varepsilon = \frac{15}{\pi^4}\frac{\sigma\varepsilon^3}{\exp(\varepsilon/T)-1^2}, \tag{1.1}$$

$$F_r = \int_0^\infty I_\varepsilon d\varepsilon,\ \int_0^\infty B_\varepsilon d\varepsilon = \sigma T^4,\ m = \rho L.$$

Here B_ε is the Planck function, ε is the photon energy, σ is the Stefan–Boltzmann constant, $\varkappa_\varepsilon$ is the spectral mass absorption coefficient, and m is the characteristic magnitude of the specific gas mass.

Let us evaluate the blackness coefficient η and the effective emission coefficient $\varkappa_e$ by means of the relationship:

$$F_r = \eta\sigma T^4 = \sigma T^4(1-\exp(-\varkappa_e m)). \tag{1.2}$$

Under the condition $\varkappa_\varepsilon m \ll 1$, which is satisfied for all radiation frequencies or photon energies ε, we obtain from (1.1)

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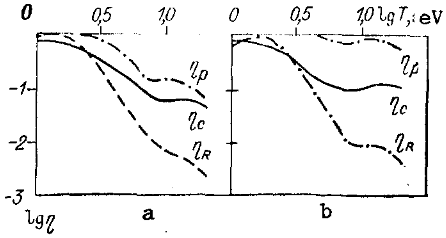

Fig. 1

$$I_\varepsilon = B_\varepsilon \varkappa_\varepsilon m,\ F_r = \sigma T^4 \varkappa_p m,\ \varkappa_e = \varkappa_p. \tag{1.3}$$

Therefore, the effective emission coefficient $\varkappa_e$ introduced goes over into the Planck coefficient $\varkappa_p$. However, the condition $\varkappa_\varepsilon m \ll 1$ is not actually satisfied in certain spectrum ranges, primarily in lines, in a sufficiently dense, multiply ionized plasma, and it is necessary to take into account reabsorption, as is done in evaluating $\varkappa_e$ by means of (1.1) and (1.2). If as a whole $\eta \ll 1$ and $\varkappa_e m \ll 1$, in the spectrum, then from (1.2)

$$F_r = \sigma T^4 \varkappa_e m,\ \operatorname{div} q_r = Q_r \rho,\ Q_r(T, \rho, m) = 4\varkappa_e \sigma T^4, \tag{1.4}$$

where Q_r is the specific power of the energy loss (per unit mass).

Computations [8] show that the dependence of $\varkappa_e$ on L or m is rather weak; correspondingly $\varkappa_e$, also turns out to be sufficiently definite if the value of L is known in advance. Taken as such in the problem of plasma heating by laser radiation is $\ell(\varepsilon_0, T, \rho)$, the path of laser radiation with photon energy ε_0.

Figure 1 shows the dependence of η on T of a bismuth plasma heated by the radiation of a neodymium laser (ε_0 = 1.16 eV) for two typical values of the bismuth plasma density $\rho = \delta\rho_L$(δ = 1; 0.1 – a and b), where the normal density of the bismuth vapors is ρ_L = 9.39 mg/cm³. The values of η_c are calculated by using (1.1) and (1.2) under the assumption that η_p and the values of $\varkappa_e = \varkappa_p$, are found when using $\varkappa_R$ as the effective mean Rosseland absorption coefficient. As is seen, the difference is sufficiently great. We note that the typical values of η for the T and δ under consideration are of the order of 0.1. Furthermore, the expression (1.4) is even extended to the case of nonuniform heating. The foundation is that the temperature of the main part of the mass of an emitting and absorbing plasma varies relatively little at exact numerical solution of the RGD problem [2-4]. The application of $\varkappa_e$ in (1.4) in every case has a much better foundation than the utilization of the Planck or Rosseland mean coefficients.

2. The system of equations describing the radially symmetric stationary motion and heating of a radiating plasma has the form

$$dp + \rho u du = 0,\ \rho u S = \dot{M},\ h = p\gamma/\rho(\gamma - 1), \tag{2.1}$$
$$\dot{M}\left[(dh + udu) + \frac{Q_r}{u}dr\right] + dF = 0.$$

Here p is the pressure, ρ the density, u the velocity, h the plasma enthalpy, γ the adiabatic index, S the area of the "channel" section in which the vapor moves ($S = 4\pi r^2$ for spherically symmetric flow), $\dot{M}$ is the mass flow rate, and F is the total energy flux of laser radiation through the section S (F = qs and q is the flux density).

The equations of laser radiation transport are

$$dF^-/dr = kF^-,\ dF^+/dr = -kF^+, \tag{2.2}$$

where F^- is the energy flux of laser radiation directed towards the target surface, F^+ is the flux of reflected radiation, and $F = F^+ - F^-$ ($F^+ > 0$, $F^- > 0$, $F < 0$), $k = k(T, \rho)$ is the linear coefficient of laser radiation absorption. On the target surface, i.e., for $r = r_0$, $F_W^+ = k_r F_W^-$ (k_r is the reflection coefficient), and phase transition conditions are also satisfied in the subsonic evaporation wave

$$\rho_w u_w S_w = \dot{M},\ p_0 = p_w + \rho_w u_w^2, \tag{2.3}$$
$$\dot{M}\left(h_w + \frac{1}{2}u_w^2\right) + F_w = -\dot{M}Q_v,\ T_w = T_v(p_w).$$

Here p_0 is the pressure ahead of the evaporation wave (in the solid), T_V is the phase transition equilibrium temperature, Q_V is the heat of evaporation, the subscript w refers to the parameters behind the evaporation wave, and F_w is the total laser and thermal radiation flux reaching the evaporation wave. Target evaporation due to the flux of thermal radiation is not taken into account below for specific computations of the problem.

On the outer boundary $r \to r_\infty$

$$F \to F_\infty = q_0 S_0,\ p \to 0. \tag{2.4}$$

The equation of state and the optical properties of the vapor are given in tabular form

$$pv = N(T,\rho) R'T,\ h = c(T,\rho) R'T = \frac{\gamma}{\gamma-1} pv, \tag{2.5}$$
$$R' = R/A,\ k = k(\varepsilon_0, T, \rho),\ \varkappa_e = \varkappa_e(\varepsilon_0, T, \rho),\ v = 1/\rho$$

(R is the universal gas constant and A is the atomic weight of the substance).

The equation of state in differential form is

$$\frac{dh}{h} - A_v \frac{dv}{v} - A_p \frac{dp}{p} = 0,\ A_p = \left(1 + \frac{\partial \ln C}{\partial \ln T}\right)\left(1 + \frac{\partial \ln N}{\partial \ln T}\right)^{-1}, \tag{2.6}$$
$$A_v = A_p\left(1 - \frac{\partial \ln N}{\partial \ln v}\right) + \frac{\partial \ln C}{\partial \ln v},\ c_s^2 = \gamma_d \frac{p}{\rho},\ \gamma_d = \frac{\gamma A_v}{1 + (A_p - 1)\gamma},\ \frac{C}{N} - 1 = \frac{1}{\gamma - 1},$$

where c_s is the speed of sound, and γ_d is the differential adiabatic index. The system (2.1) has a singular point corresponding to the passage through the speed of sound. For the constants C and N, we have $\gamma_d = \gamma$, $A_v = A_p = 1$. From (2.1) and (2.6)

$$(1 - u^2/c_s^2)\, du/u = -(dF/\dot{M}h + Q_r dr/uh + dS/S). \tag{2.7}$$

In the "sonic" section (the flow velocity u equals the speed of sound c_s), the term with du/u vanishes. We denote the parameters in this section (we call it critical) by the subscript *. If continuous gas acceleration occurs, then du does not change sign. The relationship $dF/\dot{M}h_* + Q_r^* dr/u_* h_* + dS/S_*$ should be satisfied when going through the speed of sound.

Using the transport equation (2.2), the condition $u_* = c_s^*$, and the definition of Q_r^*, we obtain in the case of a spherically symmetric flow

$$k_* r_* q_* (1 - \chi) + 2\rho_* h_* u_* = 0,\ \chi = -Q_r^* \rho_* / k_* q_* \tag{2.8}$$

(χ is the relative power of the energy losses due to radiation at the critical point).

Therefore, a definite relationship between the energy liberation power due to laser radiation (the laser radiation flux density is $q_* = F_*/S_*$, where $q_* < 0$), the energy losses by thermal radiation, and the hydrodynamic energy flux should be satisfied at the critical section.

3. To emerge from the singular point, it is necessary to execute an expansion in a small parameter. It is convenient to represent the given system and the conditions at the singular point in dimensionless form by referring all the parameters to their values in the sonic section: $\bar{r} = r/r_*$, $\bar{g} = u^2/u_*^2$, $\bar{\rho} = \rho/\rho_*$, $\bar{F} = F/F_*$, $\bar{F}^+ = F^+/F_*$, etc. For convenience, the bars over the dimensionless variables are omitted below. We write the system (2.1) and (2.2) as

$$dp = -\frac{1}{2}\gamma_d \rho\, dg,\ \rho\sqrt{g}S = 1, \tag{3.1}$$
$$dh + (\nu - 1)\, dg + \Phi_f dF + \Phi_q \frac{dr}{\sqrt{g}} Q_r = 0,\ \frac{dF^-}{dr} = k_f k F^-,\ \frac{dF^+}{dr} = -k_f k F^+.$$

Here $\Phi_q = r_* Q_r^*/h_* u_*$; $\Phi_f = F_*/\dot{M}h_*$; $\nu = 1 + \gamma_d^*/2\varphi_*$; $\varphi_* = \gamma_*/(\gamma_* - 1)$; $k_f = k_* r_*$; $\chi = \Phi_q/k_f\Phi_f$. Using (2.5) and the energy equation, we find a differential equation to determine g(r) from (3.1)

$$\frac{1}{2}\frac{dg}{dr}\frac{h}{g}\left(1 - \frac{g}{h}\Phi\right) = -\left[\frac{h}{S}2r - \frac{\Phi_f}{A_v}k_f k\,(F^+ + F^-) + \frac{\Phi_q}{A_v}\frac{Q_r}{\sqrt{g}}\right], \tag{3.2}$$
$$\Phi = \varphi\gamma_d^*/\varphi_*\gamma_d,\ \varphi = \gamma/(\gamma - 1),$$

analogous to (2.7); however, in contrast to (3.2) the real equation of state of the gas (2.5), (2.6) is used, and reflection of laser radiation from the target surface is taken into account.

At the singular point

$$r = g = p = \rho = F = S = h = k = Q_r = \Phi = 1, \quad F^+ = F_*^+/F_*,\ F^- = F_*^-/F_*,\ (F^+ + F^-) = \chi(1 + 2A_v/\Phi_q) \tag{3.3}$$

To emerge from the singular point of the system, we obtain an expansion in r at the point r = 1 to first-order accuracy, where we represent the dependences $k(T, \rho)$, $\varkappa_e(T, \rho)$, $T(h, \rho)$ near the singular point in the power-law form $k = h^{-\alpha}\rho^{\beta}$, $\varkappa_e = h^{-\alpha_e}\rho^{\beta_e}$, $T = h^{\alpha_T}\rho^{\beta_T}$, and neglect changes in A_v and Φ.

For $(r - 1) \ll 1$, $g - 1 = z_g(r - 1)$. To determine the slope z_g of the integral curve, we have a quadratic equation

$$az_g^2 + bz_g + c = 0,\ a = \frac{1}{2}\nu,\ b = -2\left[\frac{1}{2}A_v^* - (\nu - 1)(1 + \alpha) + \frac{1}{2}\beta\right] + \frac{\Phi_q}{A_v^*}\left[(\nu - 1)(\alpha + 4\alpha_T - \alpha_e) + \frac{1}{2}(4\beta_T + \beta_e - \beta) + \frac{1}{2}\right], \tag{3.4}$$

$$c = -4\left[A_v^*(1 + \alpha) + \beta - 1\right] - 2 - \frac{\Phi_f}{A_v}k_f^2 - \frac{\Phi_q}{A_v}2\left[A_v^*(\alpha + 4\alpha_T - \alpha_e) + \beta - 4\beta_T - \beta_e\right].$$

The method of solving the system (3.1) and (3.2) with the boundary conditions (2.3) and (2.4) and a singular sonic point at which the relationships (3.3) are satisfied in the case when reradiation is not taken into account ($Q_r = 0$) is elucidated in [5, 6]. We used a similar method also in the investigation of the stationary transit mode of vapor heated by a flux of fast particles or by thermal radiation fluxes in the radiant heat conduction mode [10, 11]. Consequently, we discuss here just the characteristic singularities associated with taking reradiation into account.

Integration of the system (3.1) and (3.2) starts from the singular point where the values of T_* and ρ_* are given. Values of h_*, p_*, c_{s*}, k_*, $\varkappa_e^*$, and Q_r^* are found by using tables of thermodynamic and optical properties of substances.

Integrating the energy equation (2.1) from the target surface ($r = r_0$) to the critical point ($r = r_*$) and using the boundary conditions (2.3), we find

$$h_* + \frac{u_*^2}{2} + \frac{F_*}{\dot{M}} + \frac{r_* Q_r^*}{u_*} Q = -Q_v, \tag{3.5}$$

and the dimensionless integral

$$\theta = \int_{r_0/r_*}^{1} \frac{Q_r/Q_r^*}{u/u_*} d\left(\frac{r}{r_*}\right). \tag{3.6}$$

The exact value of θ will be known only after integration of the system. Its value, from the solution without reradiation taken into account, i.e., for $Q_r^* = 0$, can be taken as the first approximation of the computations starting with modifications in which the reradiation is small ($\chi \ll 1$, $\chi = -Q_r^*\rho_*/k_*q_*$). The role of the reradiation is attenuated as T_* and r_* diminish. Later, as χ increases, extrapolation of the values of θ obtained in the previous computations can be used as the initial θ.

If the value of θ is given, then the relationship (3.5) and the conditions (3.3) permit determination of all the parameters in the critical section: r_*, F_*, and $\dot{M} = \rho_* u_* S_*$. Emergence from the singular point is realized analogously [5, 6, 10, 11] by expanding the solution near the critical section (3.4). Then the system of equations is integrated from the singular point towards the obstacle up to the intersection with the phase equilibrium curve, and r_0 and a new value of θ are determined. Substituting θ into the relation (3.5), we refine the parameters in the critical section and repeat the computation. The values of the parameters at the critical point usually converge after several iterations.

Furthermore, the computation of the supersonic flow zone is carried out (with increasing r) until F emerges at the stationary value $F = F_\infty$.

As $\chi \to 1$ the computation by the method mentioned becomes difficult, and the solution is responsive to the choice of θ. It is proposed to analyze the limit "radiation" flow mode (for $\chi \approx 1$) separately.

4. Let us estimate the plasma parameters and its thermal radiation flux under the action of a laser on a spherical bismuth target. The equation of state and optical properties of a bismuth plasma are approximated in the range T = 5-30 eV by the power-law dependencies (tables in [9] are used)

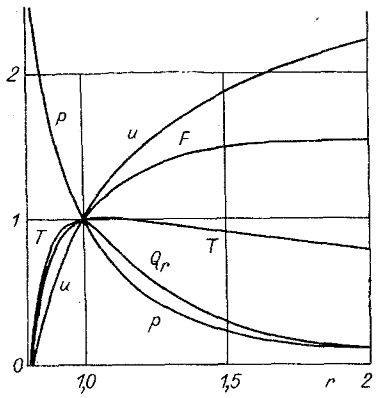

Fig. 2

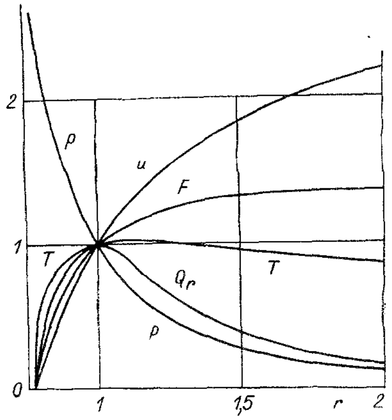

Fig. 3

$$e = 2.04T^{1.8}\delta^{-0.125},\ \varkappa_0 = 2.4\cdot 10^3 T^{1.5}\delta^{0.75}\varepsilon_0^{-2},$$
$$\eta = 0.25T^{-0.5}\delta^{-0.25}\varepsilon_0^{0.7},\ \varkappa_e = 6\cdot 10^2 T^{1.0}\delta^{0.5}\varepsilon_0^{-1.3}, \tag{4.1}$$

where e is the specific internal energy of the plasma, kJ/g, $\varkappa_0$ is the absorption coefficient of laser radiation, cm^2/g, $\varkappa_e$ is the effective absorption coefficient, cm^2/g, and $\delta = \rho/\rho_L$ is the relative density.

We estimate the plasma parameters to a first approximation by neglecting radiation energy losses ($Q_r = 0$). For sufficiently high plasma temperatures, when $h_* \gg Q_v$, the energy balance equation in the sonic section is written in the form

$$q_* = \frac{\gamma+1}{2}\rho_* u_* h_* \tag{4.2}$$

(here and henceforth $q_* > 0$). The plasma characteristic optical thickness in the critical section equals

$$k_* r_* = 1/\lambda,\ k_* = \overset{*}{\varkappa}_0 \rho_*,\ \lambda = (\gamma + 1)/4. \tag{4.3}$$

Using (4.2) and (4.3) and the approximation (4.1), we obtain relationships between the parameters in the critical section

$$T_* = 1.35 r_*^{0.23} q_*^{0.50}\varepsilon_0^{-0.46},\ \delta_* = 0.2 r_*^{-0.77} q_*^{-0.42}\varepsilon_0^{1.54}. \tag{4.4}$$

Here and henceforth, the dimensionality of the quantities is the following: q_*, MW/cm^2, r_*, cm, T_*, eV, p_*, MPa, ε_0, eV; $\delta_* = \rho_*/\rho_L$ ($\rho_L = 9.39\cdot 10^{-3}$ g/cm^3).

Using (4.4), we estimate the relative power of the radiation energy loss at the critical point

$$\chi = 4\overset{*}{\varkappa}_e \sigma T_*^4/\overset{*}{\varkappa}_0 q_* = 0.45 q_*^{0.83} r_*^{1.0}\varepsilon_0^{-1.3}. \tag{4.5}$$

The estimate (4.5) is valid in the case of low energy losses by radiation ($\chi \ll 1$); however, by using it the characteristic values of the parameters can be estimated roughly for which the radiation energy losses become significant ($\chi \approx 1$): $r_*^R = 2.2 q_*^{-0.83}\varepsilon_0^{1.3}$, $T_*^R = 1.62 q_*^{0.3}\varepsilon_0^{-0.16}$. For finite energy losses by radiation the energy balance equation is written in the form

$$q_*(1 - \chi\theta k_* r_*) = \frac{\gamma+1}{2}\rho_* h_* u_*. \tag{4.6}$$

Comparing (4.6) with (4.2), and (2.8) with (4.3), we conclude that to obtain the given values of T_* and ρ_* at the critical point when taking account of the radiation energy losses, it is necessary to increase the laser radiation flux density and to change the radius r_* or the optical thickness $\tau_* = k_* r_*$: $q_* = A_q q_*^0$, $A_q = 1 + 4\chi\theta/(1 - \chi)(\gamma + 1)$, $\tau_* = A_\tau \tau_*^0$, $A_\tau = [1 - \chi(1 - 4\theta/(\gamma + 1))]^{-1}$.

When the optical and thermodynamic properties of the vapor are given in a power-law form, the problem becomes self-similar at high temperatures ($h_* \gg Q_v$), where the self-similar profile is characterized by one parameter, for instance, the fraction χ of the energy losses by reradiation at the critical point. Here $\theta = \theta(\chi)$.

5. We shall present certain results of numerical computations of the quasistationary flow under the action of a neodymium laser on a spherical bismuth target in a vacuum. The

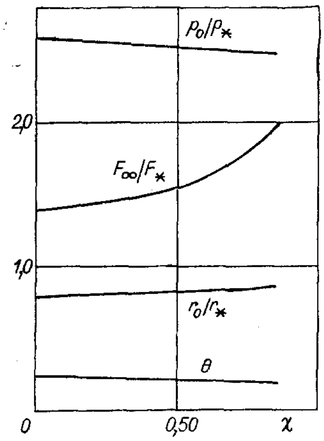

Fig. 4

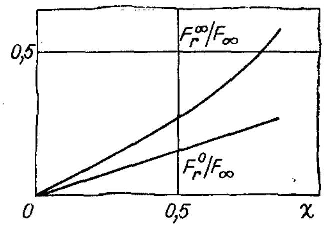

Fig. 5

parametric distributions along the radius are given in Fig. 2 for a radiation energy loss $\chi = 0.68$ at the critical point. All the quantities are referred to their values at the critical point (for $T_* = 10$ eV and $\rho^* = 2.8\cdot10^{-3}$ g/cm³, we have $r_* = 3.6\cdot10^{-2}$ cm, $F_* = 10.7$ MW, $u_* = 6.12$ km/sec, $p_* = 88$ MPa, $Q_r^* = 1.35\cdot10^{+7}$ MW/g, $q_* = F_*/4\pi r_*^2 = 0.64$ GW/cm², and the laser radiation flux density at the target in the absence of shielding is $q_0 = F_\infty/4\pi r_0^2 = 1.46$ GW/cm²).

Let us note the characteristic singularities of the flow. Near the target surface there is a heating wave front, and the maximal vapor temperature is close to the value T_* at the critical section. The radiation energy losses Q_r are maximal in the sonic section. Laser radiation energy absorption is realized at the distances $r \leqslant 2r_*$.

The flow pattern is qualitatively similar to the case when reradiation is not taken into account (compare Figs. 2 and 3 where the parameter distributions are represented for $\chi = 0$). However, there are also certain differences. Thus when the radiation energy loss is taken into account the temperature in the supersonic part of the jet, i.e., for $r > r_*$, drops somewhat more rapidly with r, while F grows more rapidly. As χ increases, the part of the laser radiation energy being absorbed in the supersonic flow zone grows, i.e., the ratio F_∞/F_* (Fig. 4) increases. The sonic point approaches the surface slightly as χ increases (the ratio r_0/r_* grows), while the ratio p_0/p_* (the ratio between the pressure at the target p_0 and the pressure p_* at the critical point) diminishes somewhat; the value θ (3.6) is reduced insignificantly as χ grows.

The thermal radiation flux incident on the target is calculated according to the relationship $F_r^0 = 2\pi \int_{r_0}^{r_\infty} (1 - \sqrt{1-(r_0/r)^2}) \times Q_r \rho r^2 dr$ in the quasivolume de-excitation approximation and under conditions of multilateral irradiation of an obstacle. The thermal radiation flux being de-excited in a vacuum is $F_r^\infty = 4\pi \int_{r_0}^{r_\infty} Q_r \rho r^2 dr - F_r^0$.

The relative fractions of de-excitation in a vacuum F_r^∞/F_∞ (F_∞ is the delivered laser radiation flux) and at the target F_r^0/F_∞ are presented in Fig. 5. For $\chi = 0.68$ $F_r^\infty/F_\infty = 0.4$, while at the target it is $F_r^0/F_\infty = 0.2$, i.e., in this case the total radiation energy losses are $\zeta \approx 60\%$. Taking account of the intrinsic plasma radiation resulted in a reduction of 1.62 times in the maximal plasma temperature T_*, while the pressure at the target was reduced 1.33 times as compared with the case when reradiation was not taken into account ($\chi = 0$). However, it should be kept in mind that obstacle evaporation by thermal radiation and its heating by the vapors being formed were not taken into account in this formulation. We note that under multilateral target irradiation conditions, additonal target evaporation owing to incident thermal radiation will result in partial compensation of the pressure diminution because of radiation energy loss. A quantitative representation of the role of this latter factor can be obtained based on results of numerical computations [2-4] of the plane problem or assumptions about absorption of the emitted radiation in an ionization and evaporation wave (the phase transition temperature of "ionization", T_V, in (2.3) is correspondingly replaced by a higher T_i at which plasma clarification occurs for this radiation).

Using the self-similarity of the problem when the power laws (4.1) are satisfied, the characteristic plasma parameters can be estimated by changing the values in the critical section while retaining the same fraction of de-excitation χ. For example, as the temperature ture diminishes to $T_* = 5$ eV for $\chi = 0.68$, $q_0 = 145$ MW/cm^2, $r_0 = 0.21$ cm, $p_0 = 40$ MPa, and $u_* = 3.3$ km/sec. The stationary mode build-up time is $t_* \approx r_*/u_* \approx 0.75$ μsec. According to experiments [12], the "scintillation" time, i.e., the time of development of a shielding plasma layer, is less than the mentioned build-up times for the stationary mode for a flux density on the order of 100 MW/cm^2.

Therefore, a quasistationary mode is established, under sufficiently prolonged exposure of a spherical target to laser radiation, for the scattering and heating of vapors intensively emitting thermal radiation here, whose role is magnified as the target size and the laser radiation flux density increase.

A number of assumptions is made for the quantitative estimate of the role of radiation effects by using the stationary model described. Thus, the phase transition is considered a quasi-equilibrium. The question of the real structure of the evaporation wave remains open. However, a comparison of computations under this assumption with experiments [12] shows that it permits a fair description of the parameters in the evaporation wave even in a mode without shielding of the plasma surface. These parameters cease to play a substantial role generally under "plasma" mode conditions, which are indeed examined here since the temperature of the main energy liberation zone exceeds the temperature in the phase transition zone greatly ($T_m \gg T_V$). We note that measurements yielded a quite fair agreement with computation by a stationary model in the range of exposure spot characteristic flux densities and radii taken as examples above. The target was aluminum here, and in conformity with estimates [6], the reradiation turned out to be an insubstantial factor. The estimates cited show that the passage over to a heavy-element target significantly magnifies the role of the radiation effects and permits their study experimentally, even under conditions similar to [13].

The action of the thermal radiation emitted by the hot plasma on the surface was neglected in the model taken above. Meanwhile, the additional target evaporation due to the thermal radiation incident on it can be compensated partially by the pressure diminution in the plasma as it cools because of radiation. On the other hand, an increase in the mass flow rate reduces the plasma temperature and attenuates the radiation. Computations were performed for a one-dimensional spherically symmetric nonstationary radiation-gasdynamic problem on the action of neodymium laser radiation on a spherical bismuth target for confirmation. Target evaporation by thermal radiation and reflection of the laser radiation from the surface (with a reflection coefficient determined experimentally in [12]) were taken into account. The flux of thermal radiation was determined by solving the radiation transport equation. The computation method is described in [14]. For a more accurate comparison with computations using this model, the plasma equation of state and its optical properties were given by the power-law dependences (4.1).

We now present the results of the computation for $q_0 = 1.46$ GW/cm^2 and $r_0 = 0.03$ cm. A computation showed that emergence in the stationary mode occurs at 0.3-0.4 μsec, which somewhat exceeds the estimate presented earlier. The parametric distribution along the radius turned out to be qualitatively quite close to that obtained above; in particular, the temperature maximum is near the sonic point. A comparatively narrow cold domain (with a temperature to 3-5 eV and 0.01 cm thickness) occurred near the target because of the additional evaporation and heating of the vapor by the thermal radiation. However, its presence did not affect the plasma parameter characteristics substantially. Thus, the temperature turned out to be only 20% lower than in the stationary model with the quasivolume de-excitation (8.1 eV instead of 10 eV), and the pressure was 10% lower (190 MPa instead of 215 MPa). The radiation losses in a vacuum were 51% (instead of 40%). All the radiation incident on the target was absorbed by it since its spectrum belongs to the far vacuum ultraviolet (20% was incident on the target in the stationary model).

Therefore, utilization of the model described above permits giving an appraisal of the role of the radiation effects (for not-too-strong de-excitation) and indicating the range of the laser pulse parameters where their experimental study and utilization are possible. A certain refinement of the stationary model can be performed by introducing the vapor evaporation and ionization waves by replacing T_v by the higher temperature T_i at which vapor bleaching occurs for the thermal radiation of the hot plasma. Further refinement requires the involvement of more tedious numerical computations (of the type described in [2-4, 14]).

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