RESONANCE INTERACTION OF A CONCENTRATED ENERGY FLUX WITH A SUBSTANCE

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Interest in the investigation of concentrated energy flux (CEF) interaction with metals is governed by the need to study the behavior and properties of metals under high energy density conditions at high temperatures and pressures, as well as by the creation of concentrated energy sources and their numerous uses in processing technology [1-4]. At this time, during consideration of the processes of CEF interaction with metals, a great deal of attention is paid to the influence of the periodicity of the action on the efficiency of materials treatment. Thus, the application of laser radiation modulation in the treatment of holes in [5] permitted substantial magnification of their depth. The presence of an optimal relative pulse duration regime when using the pulse-periodic generation mode of a CO_2—EIL (electron injection laser) for laser welding is shown experimentally in [6]. A theoretical study of the reaction of metal target evaporation on the modulation of CEF intensity [7-10] showed that the amplitude—frequency characteristics of the responses of the target surface temperature and recoil pressure amplitudes have extrema. However, the investigations in these papers were performed without taking account of energy absorption in the substance being evaporated, although energy absorption governs the dynamics of the process of CEF interaction with a substance to a great extent for a number of cases of practical importance (electron beam absorption in metal vapors [3], laser radiation absorption by a plasma being formed at the metal surface). Thus, even for a constant magnitude of the energy concentration in the CEF source—evaporated phase—condensed substance system, vibrational processes are observed [11].

Resonance properties of the beam—vapor—metal system are studied in this paper under harmonic CEF modulation on the basis of the complete system of heat and mass transfer equations with energy absorption in the evaporated phase taken into account.

1. FORMULATION OF THE PROBLEM

Let us consider the one-dimensional problem. A CEF source with energy density q_H is at the level $x = H$ (Fig. 1). In order for the boundary of the vapor—melt phase transition to remain fixed at the level $x = 0$ (in the case $q_H = \text{const}$) during evaporation, the melt is fed along the x axis at the velocity U. The value of U is determined by the magnitude of q_H, set by the gasdynamic regime in the vapor and by the material characteristics.

The system of equations describing the processes of vapor escape and energy absorption in the domain $0 \leqslant x \leqslant H$ has the form

$$\begin{gathered}\frac{\partial\rho}{\partial t}+\frac{\partial}{\partial x}(\rho u)=0,\quad \frac{\partial}{\partial t}(\rho u)+\frac{\partial}{\partial x}(p+\rho u^2)=0,\\ \frac{\partial}{\partial t}\left[\rho\left(\frac{1}{\gamma-1}\frac{p}{\rho}+\frac{u^2}{2}\right)\right]+\frac{\partial}{\partial x}\left[\rho u\left(\frac{\gamma}{\gamma-1}\frac{p}{\rho}+\frac{u^2}{2}\right)\right]-\frac{\partial q}{\partial x}=0,\\ \frac{\partial q}{\partial x}=K(\rho,T)\,q,\quad p=\frac{R}{\mu}\rho T.\end{gathered}\tag{1.1}$$

Here p, ρ, T, and u are the pressure, density, temperature, and velocity of the evaporated substance, R is the universal gas constant, μ is the molar mass, and γ is the adiabatic index. The term $\partial q/\partial x$ describes the energy absorption with absorption coefficient K(ρ, T). Specification of the absorption coefficient affords an opportunity to describe the processes in both the vapors (electron beam absorption) and the fully ionized plasma (laser radiation absorption) on the basis of the system of equations under consideration.

The heat-conduction equation for the melt is written as

$$\partial T_m/\partial t+U\partial T_m/\partial x=a_m\partial^2T_m/\partial x^2,\tag{1.2}$$

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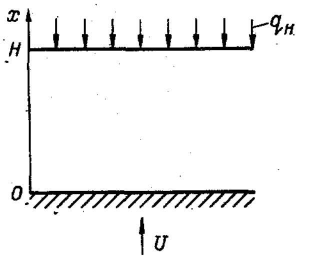

Fig. 1

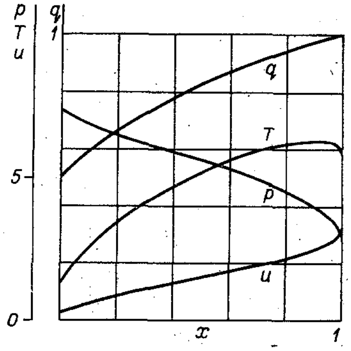

Fig. 2

where a_m is the coefficient of thermal diffusivity, and T_m is the temperature of the melt.

Equations (1.1) and (1.2) are supplemented by the following boundary conditions

$$\rho_m u_m = \rho u; \tag{1.3}$$

$$p_m = p + \rho u^2; \tag{1.4}$$

$$q = \rho u\left(L_{ev} + \frac{u^2}{2}\right) + \lambda_m \frac{\partial T_m}{\partial x}; \tag{1.5}$$

$$T_m = T; \tag{1.6}$$

$$p_m = B\exp\left(-\frac{\mu L_{ev}}{RT_m}\right); \tag{1.7}$$

$$q(H) = q_H; \tag{1.8}$$

$$u(H) = \sqrt{\gamma \frac{R}{\mu} T(H)}; \tag{1.9}$$

$$T_m(-\infty, t) = T_\infty. \tag{1.10}$$

Here B = const, L_{ev} and λ_m are the specific heat of evaporation and the heat conduction coefficient of the condensed phase, p_m, ρ_m, and u_m are the pressure, density, and velocity of this phase. Conditions (1.3)-(1.7) are written on the boundary of the phase transition, which does not agree with the level x = 0 for a time-varying quantity q_H. The first three are continuity conditions for the mass, momentum, and energy fluxes, while the relationships (1.6) and (1.7), the Clapeyron–Clausius equation, describe the equilibrium evaporation. In the general case the Jouguet condition (1.9) is, strictly speaking, not satisfied. Its validity for the stationary regime is shown in [12]. For the case of small deviations from the stationary regime, to be examined in Sec. 3, utilization of the Jouguet condition is an additional hypothesis used for the approximate analysis of the process of resonance action.

A system of equations similar to (1.1) was considered in [13]. In contrast to [13], in this paper the energy source is assumed to be at a finite distance from the target. The condition K = 0 for x > H, introduced in [12] for a long and narrow channel, corresponds physically to rapid expansion of the vapors during emergence from a channel of depth H, whereupon their density drops and they become transparent to the beam. Later hydrodynamic processes for weakly ionized vapors are considered ($K = \alpha\rho$, α = const).

2. THE STATIONARY REGIME

For a time-independent energy flux, the system (1.1) allows the stationary solution first obtained in [12]. After it has been made dimensionless, it can be written as

$$\begin{gathered} \rho u = f = \text{const}, \quad p + \rho u^2 = \varphi = \text{const}, \\ \rho u\left(\frac{\gamma}{\gamma-1}\frac{p}{\rho} + \frac{u^2}{2}\right) - Qq = \text{const}, \\ dq/dx = l\rho q, \quad p = \rho T, \end{gathered} \tag{2.1}$$

where $Q = [q]/([\rho][u]^3)$; $l = \alpha[\rho][x]$. The quantities $[q] = q_H$, $[T] = T_b$, where T_b is the material boiling point at atmospheric pressure, $[p] = 10^5$ Pa, $[\rho] = \mu[p]/(R[T])$, $[u] = (R[T]/\mu)^{1/2}$, and $[x] = H$. The dimensionless boundary conditions have the following form after a number of manipulations:

$$\varphi = \exp\{L[1 - 1/T(0)]\},$$

$$Qq(0) = f\{L + C[T(0) - \overline{T}_\infty] + u^2(0)/2\},$$
$$q(1) = 1,\ u(1) = \sqrt{\gamma T(1)}, \tag{2.2}$$

where L = μLev/(R[T]), C = μc_m/R, $\overline{T}_\infty = T_\infty/[T]$, and c_m is the specific heat of the melt. For simplicity, the energy expended in melting the solid phase is not taken into account and the specific heats of the solid and liquid metals are considered equal.

After manipulation, the system (2.1) can be reduced to one differential equation in q:

$$\frac{dq}{dx} = lfq\Bigg|\Bigg\{\frac{\gamma}{\gamma+1}\frac{\varphi}{f} \pm \sqrt{\left[\frac{\gamma}{\gamma+1}\frac{\varphi}{f} - u(0)\right]^2 - 2\frac{\gamma-1}{\gamma+1}\frac{Q}{f}[q - q(0)]}\Bigg\}. \tag{2.3}$$

The plus sign corresponds to supersonic gas flow, and the minus sign to subsonic flow. In our case (for evaporation into a vacuum), the gas velocity is less than the speed of sound [12] in the domain $0 \leqslant x < H$. Taking account of (2.1), we obtain a transcendental equation from the first two equations of (2.2) that connects T(0) with q(0) and u(0):

$$Qq(0) = \exp\{L[1 - 1/T(0)]\}\{L + C[T(0) - \overline{T}_\infty] + u^2(0)/2\}/[T(0)/u(0) + u(0)]. \tag{2.4}$$

Now, in order to solve (2.3), we proceed as follows. We give arbitrary values of q(0) and u(0) taken within reasonable limits. From (2.4) we determine T(0), and f and φ from the first two equations of (2.2). Knowing the coefficients of (2.3) and the value of q(0), this equation can be integrated numerically, and then the values of all the gasdynamic quantities can be found for x = 1. The confirmation that the two last conditions in (2.2) is satisfied permits estimating the correctness of the selection of q(0) and u(0).

The stationary distribution of parameters of the evaporated substance along x is shown in Fig. 2. The computation is executed for Q = 50 (for many metals, the values $q_H \sim 10^9$–10^{10} W/m^2 correspond to a change in Q within the limits 10–100), L = 10, l = 0.5, C = 5, $\overline{T}_\infty$ = 0.1.

3. PERIODIC ACTION

Let us consider the influence of a small periodic perturbation imposed on the constant energy flux on the gasdynamic characteristics of the vapor. We set

$$f = f_0 + f_1,\ \varphi = \varphi_0 + \varphi_1,\ \psi = \psi_0 + \psi_1,\ p = p_0 + p_1,\ \rho = \rho_0 + \rho_1,$$
$$T = T_0 + T_1,\ u = u_0 + u_1,\ q = q_0 + q_1, \tag{3.1}$$
$$\xi = 0 + \xi_1,\ T_m = T_{m0} + T_{m1}.$$

Here $f = \rho u$; $\varphi = p + \rho u^2$; $\psi = \frac{\gamma}{\gamma-1}\frac{p}{\rho} + \frac{u^2}{2}$; ξ is the coordinate of the phase transition boundary (it equals 0 for stationary evaporation). The necessity to introduce new functions f, φ, and ψ will be shown below. The quantities with subscript 0 correspond to the stationary process, the quantities with subscript 1 to small periodic perturbations having the form

$$F_1 = F_*(x)\exp(i\omega t),\quad \xi_1 = \xi_*\exp(i\omega t), \tag{3.2}$$

where F denotes each of the functions f, φ, ψ, p, ρ, T, u, q, and T_m.

Taking account of (3.1), we substitute (3.2) in the preliminary dimensionless equation (1.1). After linearization and a number of calculations, we obtain a system of ordinary differential equations in the amplitudes of the perturbations for the gasdynamic quantities

$$df_*/dx = -i\Omega\rho_*,\quad d\varphi_*/dx = -i\Omega f_*,$$
$$d\psi_*/dx = -i\Omega(\psi_*/u_0 - p_*/f_0) + \psi_0'(q_*/q_0 - u_*/u_0), \tag{3.3}$$
$$dq_*/dx = l(\rho_0 q_* + q_0\rho_*),$$

where $\Omega = \omega H/[u]$ and $[t] = \omega^{-1}$, $[f] = [\rho][u]$, $[\varphi] = [p]$, $[\psi] = [u]^2$ are taken as the scales of the quantities. The prime here and henceforth denotes differentiation with respect to x. There is a single-valued correspondence between f_*, φ_*, ψ_* and p_*, ρ_*, u_* which fails to hold only at the point x = 1. Thus

$$f_* = u_0\rho_* + \rho_0 u_*,\quad \varphi_* = p_* + u_0^2\rho_* + 2f_0 u_*,$$
$$\psi_* = \frac{\gamma}{\gamma-1}\left(\frac{p_*}{\rho_0} - \frac{p_0}{\rho_0^2}\rho_*\right) + u_0 u_*. \tag{3.4}$$

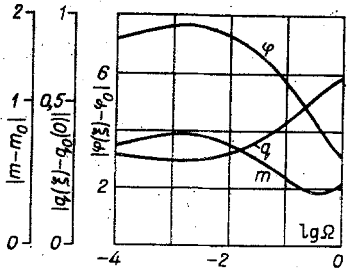

Fig. 3

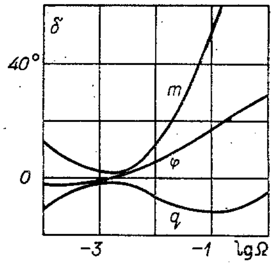

Fig. 4

Expressing p_*, ρ_*, u_* from (3.4), we obtain

$$\rho_* = \frac{-(\gamma+1)u_0 f_* + \gamma\varphi_* - (\gamma-1)\rho_0\psi_*}{\gamma T_0 - u_0^2},$$

$$u_* = \frac{(\gamma\varphi_0 f_* - \gamma f_0\varphi_*)/\rho_0^2 + (\gamma-1)u_0\psi_*}{\gamma T_0 - u_0^2}, \tag{3.5}$$

$$p_* = \frac{\left[-2\gamma u_0\varphi_0/\rho_0^2 + (\gamma+1)u_0^3\right] f_* + \left[\gamma T_0 + (\gamma-1)u_0^2\right]\varphi_* - (\gamma-1) f_0 u_0 \psi_*}{\gamma T_0 - u_0^2}.$$

If the system of equations (3.3) is written for the derivatives of p_*, ρ_*, and u_*, then the point $x = 1$ will be singular (at this point $\gamma T_0 = u_0^2$). The point $x = 1$ is not singular for the form selected for the system. The passage to the quantities f, φ, and ψ is indeed explained by this.

After making the heat conduction equation (1.2) dimensionless and substituting T_m from (3.1) with (3.2) taken into account, an equation is obtained for the amplitude of the condensed phase temperature perturbation

$$\frac{d^2T_m}{dx^2} - \frac{dT_m}{dx} - i\frac{\Omega}{k(\bar\rho f_0)^2}T_m = 0. \tag{3.6}$$

Here $k = [u]H/a_m$ and $\bar\rho = [\rho]/\rho_m$, and the quantity a_m/U is taken as the length scale. The solution of (3.6),

$$T_m(x) = C_1 e^{b_1 x} + C_2 e^{b_2 x}, \quad b_{1,2} = \frac{1}{2}\left(1 \pm \sqrt{1 + 4i\frac{\Omega}{k(\bar\rho f_0)^2}}\right)$$

should satisfy two boundary conditions: the condition of constancy of the temperature as $x \to -\infty$, and the condition of equality of the melt and vapor temperatures on the evaporation boundary. The former yields $C_2 = 0$, while the latter permits C_1 to be expressed in terms of the perturbation amplitude of the other quantities

$$C_1 = T_*(0) + \left[\frac{T_0'(0)}{k\bar\rho f_0} - T_0(0) + \bar T_\infty\right]\xi_*.$$

There are still five other conditions for the perturbation amplitudes:

the condition of substance flux continuity for $x = \xi$

$$f_*(0) = -i\frac{\Omega}{k\bar\rho^2 f_0}\xi_*; \tag{3.7}$$

the condition of energy flux continuity at $x = \xi$

$$Qq_*(0) - \left[L + u_0^2(0)/2\right]f_*(0) - f_0u_0(0)u_*(0) - Cf_0b_1C_1 + \\ \{\left[Qq_0'(0) - f_0u_0(0)u_0'(0)\right]/(k\bar\rho f_0) - Cf_0\left[T_0(0) - \bar T_\infty\right]\}\xi_* = 0; \tag{3.8}$$

the Clapeyron–Clausius equation

$$\varphi_*(0) = \frac{L\varphi_0}{T_0^2(0)}\left[T_*(0) + \frac{T_0'(0)}{k\bar\rho f_0}\xi_*\right]; \tag{3.9}$$

the condition for the energy flux q_H perturbation amplitude

$$q_*(1) = 1; \tag{3.10}$$

TABLE 1

Q	l	$k\cdot 10^{-5}$	$\Omega_{\text{res}}\cdot 10^3$	ω_{res} Hz	$q_0(H)\cdot 10^{-9}$, $\frac{\text{W}}{\text{m}^2}$	H, mm
50	0,5	2	1,8	190	3,5	6,5
50	0,1	0,4	0,85	460	3,5	1,3
10	0,5	2	0,25	27	0,7	6,5

the Jouguet condition

$$p_*(1) = p_0(1)\left[2\frac{u_*(1)}{u_0(1)} + \frac{\rho_*(1)}{\rho_0(1)}\right]. \tag{3.11}$$

The solution of the system (3.3) behaves as follows. At the point $x = 1$ $u_*(1)$ and $\rho_*(1)$ are given arbitrarily in addition to $q_*(1)$ from (3.10). The term $p_*(1)$ was determined from the condition (3.11), and $f_*(1)$, $\varphi_*(1)$, and $\psi_*(1)$ from the relationship (3.4). Such a sequence of calculations permits leaving the point $x = 1$ with the finite derivatives (3.3), which is impossible in the direct assignment of $f_*(1)$, $\varphi_*(1)$, and $\Psi_*(t)$. Determining the perturbation amplitudes at the point $x = 0$ by numerical integration of the system (3.3), the correctness of assigning $u_*(1)$ and $\rho_*(1)$ can be confirmed by satisfaction of the conditions (3.7)-(3.9).

Results of solving the system (3.3) are shown in Figs. 3 and 4 for the following values of the dimensionless numbers $Q = 50$, $L = 10$, $l = 0.5$, $C = 5$, $\overline{T}_\infty = 0.1$, $\overline{\rho} = 2.5\cdot 10^{-5}$, $k = 2\cdot 10^5$. The amplitudes of the quantities $\varphi(\xi) - \varphi_0$, $q(\xi) - q_0(0)$, $m - m_0 = \int_\xi^1 \rho dx - \int_0^1 \rho_0 dx$ are represented in Fig. 3 and the phase shift δ of these quantities relative to the phase of the quantity $q_1(1)$ in Fig. 4. The amplitude values are calculated to the accuracy of a constant that is governed by the selection of the ratio $q_*(1)/q_0(1)$. The singularity of the results obtained is the presence of a resonance regime for the pressure on the melt surface $p_m(\xi) = \varphi(\xi)$ upon energy absorption in the substance being evaporated. Let us note that according to [3, 14] the recoil pressure plays the governing role in the mechanism of "dagger" melting of the material. Results of computations for different values of Q, l, and k are presented in Table 1. If a specific metal and coefficient α are given in the energy absorption law $\partial q/\partial x = \alpha\rho q$, then it is possible to go from the dimensionless to dimensional parameters. The dimensional quantities in Table 1 are computed for iron and a 25-keV electron beam energy (here the coefficient α is inversely proportional to the square of the electron energy [3], equal to 384 m^2/kg).

An analysis of Figs. 3 and 4 shows that a minimal phase shift in the modulus of $\int_\xi^1 \rho dx$ relative to the phase $q(1)$ corresponds to its maximal value at the resonance frequency. Such agreement permits representation of the physical resonance mechanism in the system under consideration as the creation of a self-consistent regime for which a maximum mass of evaporated substance $\int_\xi^H \rho dx$ governing the energy absorption is found in the domain $\xi < x < H$ at the time of maximal energy delivery to the system. Energy absorption results in a growth of the vapor escape velocity and of the pressure in the vapor, which in the long run also produces a maximum condition for the recoil pressure on the melt surface.

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INVESTIGATION OF SLIDING SPARKS BY THE SCHLIEREN METHOD

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Sliding sparks (SS) are used for photoionization of active media of TEA CO_2 lasers, excimer lasers [1], and initiation of chemical lasers [2]. In long active laser media or at high pressures, SS are a very convenient method of preionization.

At the same time, one of the requirements which the source of preionization must meet is that the perturbation of the active medium must be minimum, especially in repetitive pulse lasers with high pulse repetition frequency. The optimal quality of the active medium is also very important when the lasing pulse duration in optical pumping of the SS is long [3].

The purpose of this work is to study the optical perturbations of the medium when SS are employed for preionization or optical pumping of the active medium in lasers.

The experiments were performed on a schlieren setup (Fig. 1). The objective lenses 3 and 4 were 30 cm in diameter and had a focal length f = 106 cm. A nonuniformity a–b was placed between the objectives. The high-quality objectives 6 and 7 with a very uniform light field, just as the remaining optical elements, were carefully aligned; otherwise it is impossible to obtain images of weak perturbations of the medium. The optimal diameter of the diaphragm was 5-2.5 mm. The ÉV-45 pulsed light source 2 (radiation duration of 250 μsec, Planck radiator at $T = 4 \cdot 10^4$ °K) transilluminated the nonuniformity under study. The schlieren motion pictures were made with the help of a fast SFR-2M camera 8. The power supply for the SS was a capacitor bank with a capacitance C = 1-3 μF, which was discharged through a commutated discharger 1 on the first winding of an IKT pulsed cable transformer with a transformer ratio of 1:10.

The SS were formed on the flat surface of a film dielectric 9, 0.8 mm thick (Fig. 1b), covering a metallic sheet 10 (initiator). The discharge gap a–b had a square-shaped protuberance and the transverse section of the spark channel was recorded on an 8-cm-long flat section of this proturberance. The discharge developed in a stable manner along the controlling strip 11; the method for controlling the development of the SS is described in detail in [4]. For monitoring the reproducibility of the conditions of discharge the voltage and current in the discharge gap were displayed on an oscillograph by a divider D and a Rogowski loop RL. The voltage oscillogram, corresponding to the photographic frames in Fig. 2b-d, is shown in Fig. 1c. The arrow marks the moment of breakdown.

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