EXPLOSIONS IN A BOUNDED VOLUME OF GAS UNDER STRONG RADIATION

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The main obstacle to utilization of strong shocks as powerful radiation sources is the shielding of their front [1]. Only quanta with energy ε less than their transparency limit ε_* can escape at large distances from the front of subcritical shocks, where this limit turns out to be somewhat below the first ionization potential I_1 of the gas over which the shock is propagated because of absorption by broadened lines in the heated layer and because of the photoeffect from the excited states of the atoms [2, 3]. Since the greatest possible values of I_1 are achieved for helium and neon and equal 24.6 and 21.6 eV, then by using the radiation of a strong shock it is not possible to produce a source in whose spectrum harder quanta would be present. For shocks with quite high velocities D that are supercritical, the hot gas in the heated layer being formed ahead of the front with a temperature close to the temperature T_S determined from the shock adiabat screens the front even in the domain $\varepsilon < \varepsilon_*$. By increasing the wave velocity D the brightness temperatures and radiation fluxes cannot successfully be raised above a specific limit thereby. Computations [2, 3] and experiment [4] show that the ultimate brightness temperature of a shock front moving in an unbounded gas is 10-12 eV for helium and neon, while the maximal radiation fluxes q_0 emitted from a front at long distances away do not exceed 0.2-0.4 GW/cm². This corresponds to an energy temperature $T_e = (q_0/\sigma)^{1/4}$ for just 7-8 eV. The ratio between the flux q_0 being emitted and the hydrodynamic energy flux $q_h = {}^1\!/_2\rho_0 D^3$ does not exceed 20%, where ρ_0 is the gas density ahead of the front.

If the size of a volume of gas over which a shock is propagated is constrained, for example, by realizing an explosion in the gas "cloud," then after the arrival of the heated layer at the boundary of such a cloud with the vacuum, the radiation fluxes will grow abruptly. The thickness of the heated layer can be great for supercritical shocks, and because of the strong radiant heat conduction an almost free yield of the radiation emitted by the front during a significant part of this time needed by the shock to traverse the whole gas volume can be realized.

This idea was confirmed in [5] by a numerical computation of the appropriate one-dimensional plane nonstationary problem of propagation of strong, intensely radiative, shock generated by a piston with constant velocity in a constrained gas layer. Computations for an almost critical shock in air showed that the radiation flux density q_0 is close to σT_S^4 during a time on the order of $\tau = L/D$, where L is the layer dimension. If heavy inert gases (xenon, krypton) or heavy metal vapors (for example, lead or bismuth) are utilized as working gases, then the shocks turn out to be supercritical for 20-70 km/sec shock velocities achievable at this time by using cumulative magnetic generators [6, 7] or explosion compressors [8, 9], and the temperature of the shock-compressed layer exceeds 15-30 eV. Still higher velocities (100-200 km/sec) can be achieved in sufficiently dense gases upon the acceleration of a thin foil by an electrical current pulse, by laser, electron, or ion beams [10].

It could be expected that the radiation flux densities would reach the value σT_S^4 upon arrival of the heated layer at the boundary with the vacuum. However, this is not so. If the mass, momentum, and energy conservation laws are examined in a quasistationary shock with the fact that part of the energy is de-excited taken into account, then the system of equations that occurs has a solution only for $0 \leqslant q_0 \leqslant q_h$. Therefore, the greatest possible energy fluxes q_0 in the quasistationary mode equal the hydrodynamic energy flux q_h. If the limit value $q_0 = q_h$ is achieved, then the parameter distribution is the following: $T = T_S$ directly at the front and all the remaining parameters also correspond to the usual shock adiabat, while we have $u = 0$, $v = 1/\rho = 0$ outside an infinitely narrow temperature peak. In order to approach the condition $q_0 = q_h$ upon the arrival of the heat layer at the boundary with the vacuum, it is necessary that 1) the heated layer be sufficiently transparent which constrains the thickness of the gas layer over which the shock moves, 2) the shock be strongly

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supercritical, i.e., $\sigma T_S^4/(1/2)\rho_0 D^3 >> 1$ would be satisfied, which imposes a definite lower bound on the shock velocity. Let us note that not only the temperature behind the front is lowered with strong de-excitation but also the magnitude of the temperature in the heated layer itself. The situation is similar to that arising in the impact of a gas jet moving at very high velocities on an obstacle [7, 11].

For an analysis of the radiation source parameters on the basis of radiation by a strong shock being formed in an explosion in a bounded gas volume that transmits radiation from the front as it is heated, it is necessary not only to consider the radiant heat transfer process but also to take into account that the shock changes its velocity as a result of an increase in the mass m_S entrained by the front and because of the energy loss by radiation.

Let us examine the development of a point explosion under the assumption of infinitely strong energy losses when compression behind the front turns out to be infinitely large. We assume that all the energy at the initial instant equals the kinetic energy of the layer, a cylindrical or spherical shell of mass m_0 with initial velocity u_0, and then deceleration of the moving gas into the working gas occurs. Because of the very strong energy losses, the thermal energy is negligible, and all the explosion energy at any time equals the kinetic energy, as before, and since the moving gas is compressed infinitely and all its particles have the identical velocity u, where u = D, then $E = (m_S + m_0)u^2/2$. The law of variation of explosion energy under the assumption of infinitely rapid de-excitation is described by the equation

$$\frac{dE}{dt} = -r_s^{j-1}\frac{1}{2}\rho_0 u^3 = -\rho_0 u \frac{E}{m_s + m_0} r_s^{j-1}. \tag{1}$$

There follows from (1)

$$dE/dm_s = -E/(m_s + m_0),\ E/E_0 = m_0/(m_0 + m_s). \tag{2}$$

If the total mass of the "working" gas equals M, then at the time of shock front arrival at the boundary $E/E_0 = m_0/(m_0 + M)$. The ratio between the de-excited and the initial energies equals $M/(m_0 + M)$. For $m_0 << M$ practically all the energy is de-excited, and the wave damping law $D \sim r_S^{-j}$ differs substantially from the analogous law $D \sim r_S^{-j/2}$ for a strong explosion with energy conservation [12], where j = 1, 2, 3 in the plane, cylindrical, and spherical cases.

The limit solutions (1), (2) are independent of the species of gas. However, the domain of applicability of the model itself depends on the optical and thermodynamic properties of the specific working gas. From this viewpoint, it is more advantageous to utilize heavy gases since they are more easily heated and permit reaching supercriticality at lower velocities. Moreover, the heated layer ahead of their front turns out to be more transparent since the radiation paths depend principally on the number of particles per unit volume [1], i.e., on the relative density δ and the quantity δ for heavy gases is smaller for equal absolute densities, and therefore, the paths are greater.

Numerical computations of the appropriate plane nonstationary radiation-gasdynamic problem were performed analogously to [5, 11, 13] to verify the reasoning described above. Since the escape of an initial dense gas layer ("explosion products") and its deceleration in a lower-density gas layer were considered, while the whole gas "cloud" was in a vacuum, then p = 0 for $m = m_0 + M$ was taken as boundary condition, where m is the Lagrange mass coordinate. Moreover, the spectral radiation intensity for the entering rays is $J_\varepsilon = 0$, and there are, correspondingly, no unilateral (integrated over the spectrum) fluxes of entering radiation: $q^- = 0$ for $m = m_0 + M$. By virtue of symmetry u = 0, $q^+ = q^-$ for m = 0. The initial data are given in the form

$$u(x) = \begin{cases} u_0 \frac{x}{L_0}, \\ 0, \end{cases} \quad \rho(x) = \begin{cases} \rho_1 = 10\rho_0 & \text{for } 0 \leqslant x \leqslant L_0 = \frac{m_0}{\rho_1}, \\ \rho_0 & \text{for } L_0 < x < L_0 + L, \end{cases} \tag{3}$$
$$e(x) = 0,\ 0 \leqslant x \leqslant L_0 + L,$$

TABLE 1

lg δ	v_0, km/sec	L, cm	T_0, eV	Q_0, GW/cm²	t_L, μsec	E_0, kJ/g
−0,5	95	0,34	64	160	0,036	0,56
−1	64	2,0	35	15	0,31	0,47
−1,5	43	6,0	20	1,5	1,4	0,20

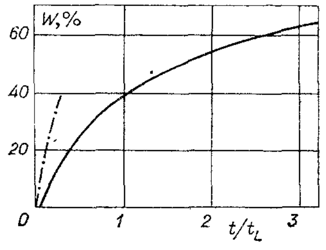

Fig. 1

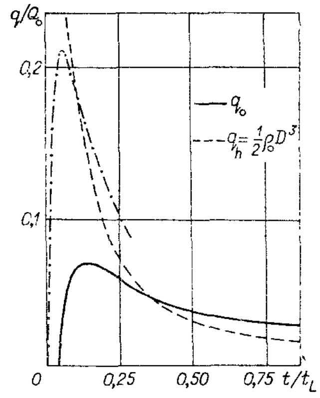

Fig. 2

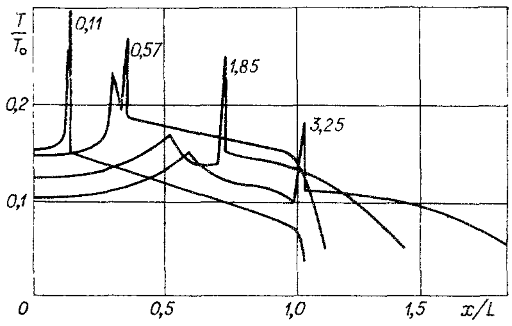

Fig. 3

where $L = M/\rho_0$ is the initial thickness of the decelerating layer. The velocity and temperature distributions mentioned of the "products" correspond to their inertial escape prior to the beginning of deceleration. Xenon was taken as working gas. To convert the results for the different density of xenon, the equation of state and the dependence of the Rosseland path l_R on T and δ [14] were approximated by power laws [8] in the range to T = 30 eV:

$$T = 0.42\, e^{0.61}\delta^{0.09},\quad l_R = 0.0017\, e^{0.92}\delta^{-1.57}, \tag{4}$$

where T is the temperature, eV; e is the internal energy, kJ/g; l_R is the Rosseland path, cm. The r.m.s. velocity of the decelerating gas equals $v_0 = (1/\sqrt{3})u_0 = 0.58u_0$. The following characteristic quantities

$$v_0, L, E_0, T_0 = T\left(e = \frac{v_0^2}{2}, \delta\right), \; Q_0 = \rho_0 v_0^3, \; t_L = \frac{L}{v_0}, E_0 = \frac{\rho_0 L_0 v_0^2}{6}.$$

are presented for different densities in the table for the case $L = 50L_0$ (for the gas mass $M = \rho_0 L = 5m_0$).

The dependence of the de-excitation W on the time is presented in Fig. 1. It is seen that up to the time of shock front arrival at the boundary with the vacuum, 65% of the initial energy E_0 is de-excited (by estimates (2) it should have been 75%). Time dependences of the emerging radiation flux density q_0 and the hydrodynamic energy flux q_h are presented in Fig. 2. The maximum q_0 is 12, 1.1, and 0.11 GW/cm² for δ = 0.3162, 0.1, and 0.0362 (the maximal energetic temperatures $T_e^m = (q_0^m/\sigma)^{1/4}$ are 18, 10, and 5.6 eV, respectively).

For $t/t_L < 0.5$ the wave is strongly supercritical ($\sigma T_s^4/\frac{1}{2}\rho_0 D^3 \sim 5$–$10$), hence, $q_0 \simeq q_h$. For $t/t_L > 0.5$ the quantities T_- and T_s are commensurate and q_0 exceeds q_h almost twice. Results of an analogous computation with the same value of L_0 but a smaller value of L, namely, $L = 10L_0$, are represented in Figs. 1 and 2 by dash-dot lines (in this case $M = m_0$). Since the dimension L is less here, the heated layer emerges differently on the boundary with the vacuum, and the wave velocity at the time of arrival is greater and the maximum q_0 is correspondingly greater (approximately triple). However, the relative de-excitation here turns out to be halved. As the layer mass in which the gas is decelerated diminishes, the value of the radiation, energetic and brightness temperature fluxes can be increased still more for the same shock velocity.

Presented in Fig. 3 are the temperature distributions at different times t (the values of t/t_L are indicated at the appropriate curves) for the version with $L = 50L_0$. As is seen the pattern of the process obtained in the numerical computation corresponds qualitatively to that described above.

The estimates and computations presented show that more than 50% of the explosion energy is de-excited upon the shock arrival at the boundary with the vacuum (almost all the energy in the limit, for very strong shocks). The radiation flux density will here be determined mainly by the hydrodynamic energy flux $\frac{1}{2}\rho_0 D^3$. The limit pattern is independent of the thermodynamic and optical properties of the working gas. However, the approximation to the limit pattern is facilitated when utilizing heavy gases.

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STRUCTURE OF SHOCK-WAVE FLOWS WITH PHASE TRANSITIONS IN IRON NEAR A FREE SURFACE

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A compression shock wave with a stable three-front configuration, associated with a polymorphic phase transition, was observed in Armco iron in [1]. The $\alpha \rightleftarrows \varepsilon$ transformation in iron was carefully studied in static tests in [2], which discovered the martensitic character of the phase transition and showed that the α- and ε-phases of iron coexist within a broad range of pressures corresponding to the beginning of the forward and reverse transitions. In [3] manganin pressure transducers were used to record directly the multifront structure of both a compressive shock wave and a rarefaction wave at internal points of an iron specimen. Shock unloading waves were also recorded experimentally. In [4] laser interferometry was used to determine accurately the velocity profile of a free surface of a shock-loaded iron target. The most complete study of the polymorphic transformation in iron under dynamic loading conditions was made in [5], where again laser interferometry was used to obtain detailed measurements of the velocity of the free surface of Armco iron specimens loaded by shock waves of different intensities. The investigation uncovered fine-scale shock-wave effects connected with the arrival of a three-front shock wave at the free surface. In particular, it was shown that under certain conditions an additional fourth velocity jump occurs in the velocity profile of the free surface behind the third wave. The experiments were conducted for specimens of different thickness within the stress range from 10 to 40 GPa.

Numerical studies connected with the travel of shock waves in solids and the occurrence of associated physicochemical effects were performed in [6, 7], which developed a model of an elastic-plastic body with phase transitions and proposed phase-transition kinetics. The

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