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STRUCTURE AND BRIGHTNESS OF NONSTEADY-STATE SUPERCRITICAL SHOCK WAVES IN AIR OF REDUCED DENSITY

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Basic ideas about the quasisteady structure of intensely emitting shock waves (SW) which were given in [1-3] have made it possible to make qualitative estimates of the thickness of a heated layer and peak behind the SW front approximating a gray gas, and also to estimate SW brightness at different velocities in air of normal density [3-5]. Numerical calculations for the nonsteady-state radiation-gas dynamic problem of SW propagation in air with velocities up to 50-70 km/sec have been carried out for densities of 0.1-0.03 of normal density [6-9]. The emission spectrum of an air plasma (up to 500 spectral intervals have been introduced in order to describe it) has been considered in detail, and detailed tables of thermodynamic [10] and optical [11] properties of air have been used. In order to reduce the volume of computations a special method has been used for averaging equations of emission transfer [12, 13]. Calculations have made it possible to determine the dependence of the brightness temperature of emission on its wavelength with different wave velocities and air densities, and also the time for emergence of parameters into their steady-state values.

Numerical calculations [6-9] have confirmed the validity of the basic qualitative ideas [1-5], and made it possible to define quantitative characteristics for emitting SW. However, new qualitative features of these SW were detected, i.e., occurrence of a two-region structure for the heated layer ahead of the wave front. Analysis of emission spectra and the nature of change in group and integral (according to spectrum) one-sided emission flows has shown that the reason for occurrence of a two-region structure is the difference in behavior of absorption coefficients in different parts of the spectrum with a change in temperature. The drop in absorption coefficient for fronts with energies of 6.5-11 eV at temperatures of 0.7-0.9 eV, connected with dissociation of the air molecules, leads to occurrence of a heating wave and brightening. High energy quanta, which have a capacity for photo-ionization, form a hot region adjacent to the SW front. Between the ionization and dissociation waves an extended, comparatively cold zone arises whose existence has not previously been noted.

The role of radiation processes may be characterized by the parameter $\chi = q_r/q_h$, where q_r is radiation flux density; q_h is hydrodynamic radiation flux $\left(q_h = \frac{1}{2}\rho_0 D_s u_s^2\right)$, ρ_0 is gas density ahead of the SW front; D_s is SW velocity; u_s is gas velocity behind the SW front. If the plasma behind the SW front is optically thick and uniformly heated, then $q_r = q_b = \sigma T_s^4$

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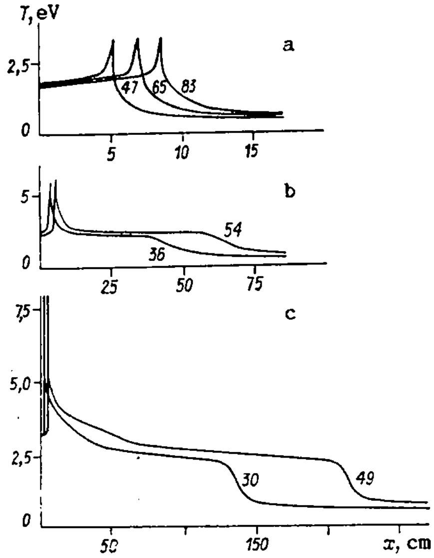

Fig. 1

(q_b is radiation flux of a black body at temperature T_s according to the Hugoniot adiabat; σ is Stefan—Boltzmann constant). We approximate the equation of state [10] by an analytical relationship in the temperature range 0.8-30 eV.

$$e = 31T^{1.5}\delta^{-0.12}. \tag{1}$$

Here e is the amount of internal energy for a unit mass of gas in kJ/g; T is temperature in eV; δ is ratio of density ρ for air to its normal density ρ_L = 0.00129 g/cm³. We write normal relationships at an intense SW front

$$D_s = \frac{\gamma_s + 1}{2} u_s, \quad e_s = \frac{u_s^2}{2}, \quad \theta_s = \frac{\rho_s}{\rho_0} = \frac{\gamma_s + 1}{\gamma_s - 1}, \quad p_s = \rho_0 D_s u_s, \tag{2}$$

where p_s is pressure; index s relates to parameters behind the SW front. By assuming the effective index for the adiabat γ in the range of temperature and density in question is constant and equal to 1.2, we obtain

$$T_s = 0.077 u_s^{1.33} \delta_0^{0.08}, \quad q_h = 0.071 u_s^3 \delta_0, \quad p_s = 1.4 u_s^2 \delta_0,$$
$$q_b = \sigma T_s^4 = 3.6 \cdot 10^{-6} u_s^{5.32} \delta_0^{0.32}, \quad \chi_s = 0.52 \cdot 10^{-4} u_s^{2.32} \delta_0^{-0.68}. \tag{3}$$

Here D_s and u_s are in km/sec; p_s is in MPa; T_s is in eV; q_h and q_b are in MW/cm².

In those cases when the plasma behind the SW front does not have an optical thickness, for characteristics of the emitting capacities of the plasma an idea was introduced in [13] of an effective average absorption coefficient k_e determined as follows. We calculate the integral (over the whole spectrum) intensity G of a uniformly heated layer of temperature T, prescribed density ρ and thickness x, and by using real values of spectral absorption coefficients [11] $k_\varepsilon(\varepsilon, T, \rho)$:

$$G(T, \rho, x) = \int_0^\infty B_\varepsilon [1 - \exp(-k_\varepsilon x)]\, d\varepsilon, \quad B_\varepsilon = \frac{15}{\pi^4} \frac{\sigma \varepsilon^3}{\exp(\varepsilon/T) - 1} \tag{4}$$

(ε is proton energy; B_ε is Planck function). The value of $k_e(T, \rho, x)$ should give the same value of $G(T, \rho, x)$, as the solution of the spectral problem (4):

$$G(T, \rho, x) = \sigma T^4 [1 - \exp(-k_e x)]. \tag{5}$$

With $k_e \ll 1$ from (5) an expression follows similar to volumetric brightening $G = \sigma T^4 k_e x$. However, strictly speaking the value of k_e coincides with the Planck absorption coefficient k_p only for plasma which is so thin that $k_\varepsilon x \ll 1$ for all ε. A relationship $k_e(T)$ was given in [14] for different pressures p and thicknesses x, and in [13] the dependence was given for k_e on x for different T and p. These values in the range T = 1.2-4 eV, p = 0.3-10 MPa, x = 1-100 cm can be approximated by the expression $l_e = 7.0\sqrt{x}/p$. Here $l_e = 1/k_e$ is in cm.

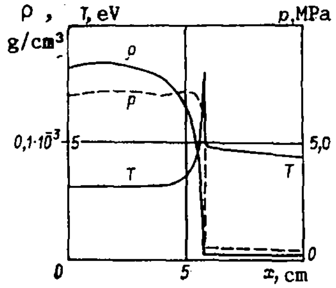

Fig. 2

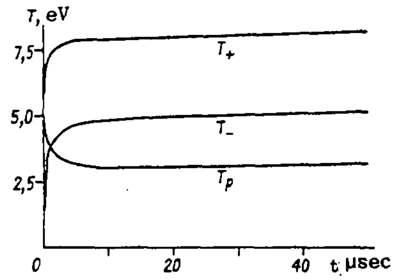

Fig. 3

We consider SW propagation with $\delta = 1/300$. With a piston velocity of 20 km/sec the wave is already supercritical ($\chi = 2.2$). Temperature behind the front (according to calculation from the shock adiabat) $T_S = 2.7$ eV. Characteristic energy for the emitted radiation is 7-8 eV. Pressure behind the front $p_S = 1.9$ MPa, and therefore darkness for the shock-compressed layer sets in with a thickness of the order of 10 cm, i.e., with distance x_S covered by the front of the order of 1 m.

With a temperature behind the front of 2-3 eV the maximum of the black body spectrum lies in the region 6-9 eV, i.e., in just that region which provides advance of the dissociation wave. Under quasisteady conditions due to absorption this zone stops at a certain distance from the front. Loss of energy "at infinity" is small since emitted and absorbed energy is taken into the SW front afresh together with heated gas. However, the time for establishing this structure may be very large and correspondingly the distance covered by the front before establishing this structure is large. It may be much greater than the time when nontransparency conditions are obtained behind the front. If it is assumed that all of the emitted energy with a flow density approximately equal to hydrodynamic flow, i.e., 1.8 MW/cm^2, is consumed in forming the heated layer with a temperature of 0.9 eV or an internal energy of about 40 kJ/g or 0.15 J/cm^3, then dissociation wave velocity with respect to the piston generating the SW is about 120 km/sec. Paths of emission l_e with photon energies $\varepsilon \sim 6$-9 eV at $T < 0.8$-0.9 eV exceed 30 m, and therefore the time for establishing a steady-state structure is more than 300 μsec.

Since the distance between the SW front and the ionization wave increases continuously, the relationships of a steady-state balance usually used are not applicable. It is possible to assume that all of the emitted energy emerges almost freely from the front and for the zone behind the front gas cooling is significant. It reduces the average energy of the emitted quanta, and even more strongly it increases the time for establishing a quasisteady-state regime.

Now we consider a high gas velocity behind the SW front ($u_S = 40$ km/sec). According to relationship (3) the temperature behind the front is 6.6 eV and the maximum emission spectrum emitted by a dark body with this temperature is in the quantum region of 18 eV, i.e., in the region between the first and second oxygen and hydrogen ionization potentials. A SW with this velocity is strongly supercritical ($\chi_S = 13$). According to classical theory [1-3] energy transfer in the heated layer depends on the mechanism of radiant thermal conductivity, and temperature distribution has the nature of a nonlinear thermal wave with a maximum temperature equal to T_S:

$$\frac{T}{T_s} = \left(1 - \frac{x}{x_T}\right)^{1/\omega}, \quad x_T = \frac{16\chi_s l_R^s}{3\omega}, \quad l_R^s = l_R(T_s, \rho_0), \quad \omega = 4 + a - b.$$

Here x_T is heated zone thickness; l_R is the average for emission travel; a and b are exponents in the approximation rules $l_R(T, \rho) = BT^b f_1(\rho)$, $e(T, \rho) = AT^a f_2(\rho)$. According to [11] values of l_R in the range T = 3-10 eV are approximately constant and with $\delta = 3\cdot10^{-3}$ they equal 10 m. Since b = 0, whereas according to relationship $a = 1.5$, then $\omega = 5.5$ and $x_T =$ 120 m. Thus, the thickness of the heated layer in the steady stage is large and considerable time is necessary in order to form a quasisteady-state structure. In the nonsteady-state the maximum temperature ahead of the front will be much less than T_S. Travel of emission in cold air for quanta with energies of 18-30 eV are $2\cdot10^{-3}\ \delta^{-1}$, and with $\delta = 3\cdot10^{-3}$ they reach in all only 0.6 cm. Therefore, energy emitted by the SW front and absorbed ahead of it (due to the photoelectric effect) rapidly leads to an increase in temperature compared

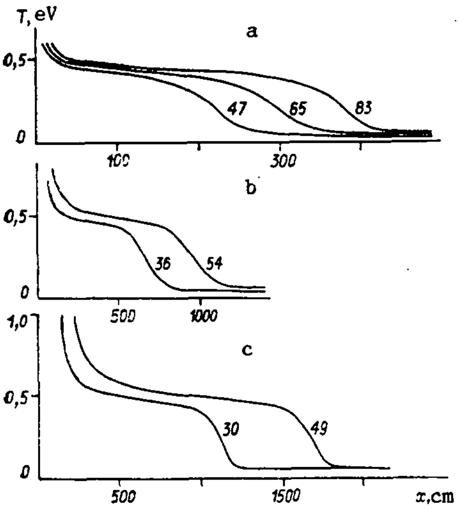

Fig. 4

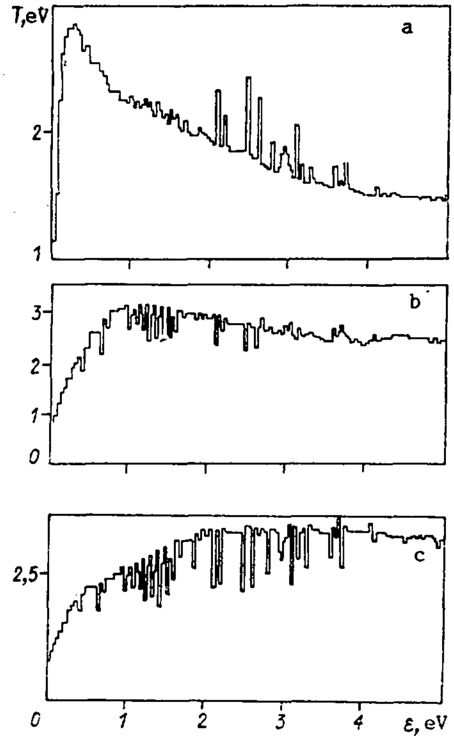

Fig. 5

with the initial value. However, with approximation of temperature to about 3 eV single-stage ionization of the air plasma and marked brightening of it sets in, which provides occurrence and advance of an ionization wave. With T = 3 eV and $\delta = 3 \cdot 10^{-3}$ the internal energy of a unit mass e = 300 kJ/g, and of a unit volume $e\rho = 1.4$ J/cm^3. With a radiation flux of 27 MW/cm^2 we find the ionization wave velocity in relation to the emitting front is 190 km/sec, and the ionization wave travels a distance of the order of 10 m in 50 μsec. As a result of intense radiation the gas temperature behind the front appears to be much lower than T_s. We estimate the value of this temperature by proceeding from the condition that the radiation flux σT_r^4 with this temperature T_r equals the hydrodynamic flow through the front q_h. Then from (1) we obtain $T_r = 0.91 u_s^{0.75}\delta_0^{0.25}$, $\frac{T_r}{T_s} = 12 u_s^{-0.58}\delta_0^{0.17} = \left(\frac{u_s}{u_s^k}\right)^{-0.58}$. Here u_s^k is critical SW velocity. With $u_s = 40$ km/sec, $T_r = 3.5$ eV, $T_r/T_s = 0.53$.

Thus, the estimates made show that an intermediate quasisteady-state stage exists for propagation of powerful intensely emitting SW when on one hand the plasma behind the SW front became nontransparent, and on the other hand complete steadiness was still far from achieved. In a rarefied plasma this stage may be very prolonged. For supercritical waves in this intermediate stage the temperature ahead of the front appears much lower than follows from classical theory, i.e., much less than the temperature from the shock adiabat. Behind the wave front plasma cooling is very marked.

The estimates made need refining since temperature distribution behind the front is far from uniform and strictly speaking the wave structure is not steady. Therefore, we carried out numerical solutions for the set of unidimensional nonsteady-state gas dynamic equations and spectral equations for transfer of emission. The calculation procedure is entirely similar to that used in [6-9] and therefore it is not described here. We give some results of calculations for $\delta = 1/300$ and velocity u_p of the piston generating the SW equal to 20, 30, and 40 km/sec.

Presented in Fig. 1a-c are distributions of temperature T over distance x read from the piston at different instants of time t (values of t in μsec are shown on the curves) for piston velocities u_p = 20, 30, and 40 km/sec. Clearly seen in Fig. 1b, c is ionization wave propagation with a temperature behind the front of about of 2 eV. The width of the front is about 15-20 cm. Ahead of the hydrodynamic jump there is an increase in temperature connected with absorption of the hard part of the emission spectrum. Behind the SW front a temperature peak is seen. Maximum temperature T_+ is achieved directly behind the hydrodynamic jump propagating through heated gas. The drop in temperature in the region of the peak is connected with intense cooling as a result of radiation.

In Fig. 2 for instant of time 50 μsec the distribution of temperature T(x) is shown on a large scale in shock-compressed and in heated layers close to the front for u_p = 40 km/sec. It can be seen that the temperature close to the piston is much lower than the temperature ahead of the front and it is about 3 eV, whereas ahead of the front it is close to 5 eV, that is to say a different value much lower than temperature T_s from the shock adiabat. Also given is the distribution of pressure p and density ρ. The maximum density is $0.17 \cdot 10^{-3}$ g/cm^3, i.e., the maximum compression reaches 40 which far exceeds the normal compression behind the front of intense SW calculated from the shock adiabat [10]. Pressure p in the shock-compressed layer is almost unchanged and it remains close to the value from the shock adiabat p_s. An increase in compression is connected with intense gas cooling as a result of radiation in the region ahead of the front (the piston was assumed to be thermally insulated). It is noted that at the instant of time in question the shock-compressed layer is not transparent for natural emission, the average travel with a temperature of 3 eV and pressure 7 MPa is about 2 cm, whereas the thickness of the shock-compressed layer is about 6 cm. Estimates show that in this stage the effect of actual piston characteristics ceases to show.

Given in Fig. 3 for u_p = 40 km/sec is the dependence on time t of values T_p, T_-, and T_+. As can be seen, starting from time 10-20 μsec quasisteady-state distributions for temperature behind the front and close to it are established, the values of T_+, T_-, and T_p differ markedly from values prescribed by classical theory for intensely emitting SW [1-5]. A similar picture also occurs for a velocity of 30 km/sec, and the time for establishment is about the same as for 40 km/sec, but for a velocity of 20 km/sec it is markedly greater (about 50-60 μsec).

Presented in Fig. 4a-c is temperature distribution ahead of the front in the so-called cold zone for velocities of 20, 30, and 40 km/sec respectively. Dissociation wave propagation is clearly seen with a temperature behind the front of 0.4-0.5 eV. The width of the dissociation zone markedly exceeds the thickness of the ionization region in the same instants of time. For u_p = 20 km/sec the velocity of the leading front of the wave in question is almost unchanged with time up to 80 μsec when it passes 400 cm (average propagation velocity in relation to the piston 50 km/sec). For u_p = 40 km/sec this wave travels 1700 cm at instant of time 50 μsec (average velocity 340 km/sec). The maximum velocity in this case at the start of the process was about 500 km/sec, and at instant of time 50 μsec it fell to about 230 km/sec which is connected with the start of noticeable absorption of emission in the region ahead of the dissociation wave front. For a velocity of 30 km/sec the dissociation wave draws away from the piston by 1000 cm at instant 55 μsec (average velocity 200 km/sec) and the drop in velocity is still small (it does not exceed 30% compared with the maximum value). Thus, calculations have confirmed the existence of an intermediate stage in which a quasisteady-state structure is established behind the SW front and close to it, and nonsteady-state dissociation and ionization waves move in front.

Shown in Fig. 5a-c is an emission spectrum, more accurately the dependence of brightness temperature T_ε on photon energy ε for u_p = 20, 30, and 40 km/sec at instants 83, 55, and 50 μsec respectively, i.e., in the intermediate quasisteady-state indicated. For u_p = 20 km/sec, T_ε is greatest in the IR-region of the spectrum. It is governed by emission at the peak (T_+ = 3.5 eV). At the same time in the visible and near UV-regions (ε = 3-5 eV) it appears markedly lower than even T_p (by about 1.4 eV).

For u_p = 40 km/sec the maximum for the spectrum lies in the visible region and it is 3.2 eV, which is close to T_p and much lower than T_s, and moreso T_+ at the peak. In this case emission at the peak is absorbed close to the front, but at considerable distances it leaves emission from the colder regions. For u_p = 30 km/sec the value of T_ε is almost constant and it is 2.5-3 eV, which is somewhat higher than T_p = 2.4 eV. The density of the radiation flux moving at considerable distances from the front is about 0.5, 1.5, and 2.0 MW/cm^2 for u_p = 20, 30, and 40 km/sec, which is 28, 25, and 14% of the hydrodynamic flow q_h. Even greater radiation flux may be obtained within a cold heated layer of considerable extent (at the level of about 0.7, 2.0, and 3.0 MW/cm^2 respectively). Thus, supercritical SW may appear to be quite intense sources of thermal emission (with an efficiency of the order of 15-30%).

Now we consider some possible applications of the results obtained. With movement of large meteoric bodies in the Earth's atmosphere with very high velocities ahead of the leading SW it is possible for heated layers to occur. The characteristics of these bodies with steady-state flow around the body may be estimated on the basis of results for solving the nonsteady-state problem for movement of the SW generated by a piston [6], and the results should be used for calculations carried out with distances travelled by the piston of the order of the size of the body. The dimension of these large meteoric bodies may reach meters, tens, or even hundreds of meters [15, 16]. As can be seen from the estimates made, for a height of the order of 40-50 km or more a classical steady-state SW structure is not established and estimates of their parameters should be carried out on the basis of more accurate solutions taking account in particular of the possibility of some intermediate quasisteady-state structure including that considered.

In conclusion we note that this structure may arise not only for intense SW in air, but also in other gases, not necessarily molecular gases. This confirms direct calculations carried out for example for xenon with reduced (compared with normal) density. The qualitative picture of SW propagation in xenon does not differ from that described above for air. However, supercritical wave amplitudes are obtained with low velocities and sizes of devices, which makes it possible to carry experimental verification of the results obtained. It is noted that in xenon the temperature ahead of the front T_- appeared to be higher than that close to the piston T_p. This fact, which contradicts classical theory [1-5] (for steady-state SW), is explained by the marked nonuniformity of emission close to the hydrodynamic jump, which is a spectral effect.

The approximate local thermodynamic equilibrium used in the calculations could be disturbed in a narrow temperature peak of the front, which could lead to a marked reduction in the maximum of electron temperature [17]. However, estimates according to [5] show that in the case in question the time of electron-ionic relaxation is 10^{-9}-10^{-8} sec, and the residence time for particles in the peak is 10^{-7} sec or more. Therefore, the role of nonequilibrium is apparently small.

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COLLAPSE OF A SPHERICAL CAVITY IN A MEDIUM COMPLETELY TRANSPARENT TO VOLUME RADIATION

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When a spherical cavity collapses there develop in the vicinity of the center a number of unique features in flow gas dynamic characteristics, which can essentially be described by a self-similar solution corresponding to the process under consideration.

Self-similar solutions for gas dynamic flows upon collapse of a spherical cavity have been found using the assumption of flow isoentropicity [1]. In the present study the flow will be considered with presence of radiation losses in the medium, which is completely transparent to volume radiation, these losses developing when the gas temperature outside the cavity is sufficiently high. It will be assumed that the character of the radiation corresponds to a braking mechanism of free—free electron transitions, since at sufficiently high temperature all atoms of the material are completely ionized. The gas dynamic equations, the self-similar solution of which will be obtained below, differ from the classical system only in the presence of a term corresponding to radiant losses in the energy equation [2]:

$$\frac{\partial \rho\left(e+\frac{u^2}{2}\right)r^2}{\partial t}+\frac{\partial \rho u\left(e+\frac{u^2}{2}+\frac{p}{\rho}\right)r^2}{\partial r}=Q_0 r^2\rho^\alpha T^\beta$$

(where the constant $Q_0 < 0$, $\alpha = 2$, $\beta = 1/2$). Nevertheless this addition significantly changes the character of the flow: it becomes nonisoentropic; the self-similarity index increases as compared to the index obtained without consideration of radiant losses, which intensifies cumulation. The self-similar solution will be determined using the principles of [1], but will be significantly more complicated in view of the absence of an adiabatic interval in the case under consideration.

1. Mathematical Formulation of the Problem. In the self-similar solution the equation of state for the gas is assumed polytropic:

$$p=\rho c^2/\varkappa,\ e=p/[(\varkappa-1)\rho],\ T=p/(R\rho)=c^2/(\varkappa R).$$

Here p is pressure; e is specific internal energy; ρ is density; T, temperature; c, speed of sound; R, universal gas constant; $\varkappa$, polytropy index.

For collapse of a spherical cavity r and t are the distance from the center of the cavity and the time measured from the moment of collapse ($t < 0$). After transition to dimensionless variables

$$t=t_0t,\quad r=r_0r,\quad c^2=r_0^2c^2/t_0^2,\quad u=r_0u/t_0,\quad \rho=\rho_0\rho,\quad p=\rho_0r_0^2p/t_0^2,$$

where any two quantities (for example, ρ_0 and t_0) are arbitrary positive constants with appropriate dimensions, and

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