NUMERICAL CALCULATION OF THE PROPAGATION OF A PLANE SUBSONIC RADIATION WAVE THROUGH A GAS IN OPPOSITION TO A FLOW OF LIGHT RADIATION

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The propagation of a plane heating and ionization wave through a gas is considered; the wave is sustained by a strong flow of monochromatic optical radiation (traveling in the opposite direction) through energy transfer attributable to the emission of a continuous spectrum. In the range of radiation flux densities under consideration, a situation arises in which the expanding hot layer generates a shock wave transparent to the incident radiation. The radiation wave is subsonic. The pressure within the hot layer is smoothly distributed, so that its parameters may be determined by considering the equations of energy and transport of the monochromatic source radiation and the radiative-transfer equations for various frequencies and directions. The true spectral composition and distribution of the radiation are considered in detail, using refined tables of the thermodynamic and optical properties. The results of numerical calculations relating to air are presented; so are certain details of the methods used in averaging the transfer equations, which prove very efficient for the radiation-gas-dynamic problem under consideration and greatly reduce the volume of calculations.

1. Powerful sources of monochromatic optical radiation are widely employed for heating gases to high temperatures. In addition to the question of achieving thermonuclear temperatures of the order of 1-10 keV [1], which requires very high radiation flux densities there is also the problem of heating plasma to lower temperatures such as 1-10 eV as required in various fields of technology. The production of plasma with temperatures of this order under laboratory conditions facilitates a number of physical and gasdynamic investigations, in particular those concerned with determining the optical and thermodynamic properties of gases. It is especially interesting to determine such properties for dense plasma. In addition to this the heated gas emits strong radiation, the long-wave part of this radiation being irrecoverably lost from the plasma "cloud." This makes it possible to create strong sources of radiation with a continuous spectrum.

The question arises as to the relationship between the temperature achieved and the density of the radiation incident upon the plasma space, the temperature and density distribution within this volume, the rates of flow, and the spectrum of the radiation generated in the plasma which passes outside the boundaries of the heated volume, i.e., the question of the "optical plasmotron" characteristics.

The achievement of high densities in the plasma is eased when the plasma cloud is surrounded by a fairly dense gas which restricts disintegration ("high-pressure plasmotron"). The simplest conditions for carrying out physical investigations and also for theoretically predicting the parameters and subsequently comparing them with experiment arise when the motion and transfer of the energy take place under plane conditions. It is the plane case which we shall be considering in this paper.

Let us assume that a plane ionized layer capable of absorbing the optical radiation of the external source very strongly is first created in the gas. Let us consider the evolution of the parameters in this layer.

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One of the methods of creating such a layer lies in the action of radiation on the surface of a solid target placed in the gas. After a certain holding period associated with the transfer of heat into the interior of the surface, intensive evaporation begins, and if the radiation flux density exceeds a certain critical value a "flash of absorption" arises in the layer of vapor [2-4]; evaporation ceases, since the radiation is practically all absorbed in the ionized layer so formed.

Under conditions in which the disintegration of the vapor is restricted by the gas medium, the absorbing layer heats rapidly; it starts radiating very strongly and heats the layers of gas at the interface with the vapor. Subsequently the process develops in the gas surrounding the target. Calculations based on the method of [2-4] for the range of flux densities here considered lead to comparatively short times of development of the absorbing layer. Thus for a flux density of the order of 10 MW/cm^2 and an aluminum target this time is 15 μsec. The flash time may be still further shortened by using a brief pulse of initiating radiation. Since the screening time (according to [2-4]) diminishes rapidly with increasing flux density, the energy of the "igniting pulse" may be extremely low. Bearing in mind this mode of initiation we shall assume that the ionized layer of vapor lies above a stationary plane surface.

For a sufficiently high temperature the main part of the radiation emitted by the plasma lies in the ultraviolet range, a considerable proportion of this radiation belonging to quanta with energies greater than the first ionization potential. This radiation is trapped in the cold layers of gas adjacent to the plasma cloud and heats them. The latter in turn start absorbing the optical radiation of the source very strongly; they heat rapidly to a temperature close to the maximum, and themselves start emitting strong radiation, heating the following layers, and so on. A wave of absorption and heating starts moving through the cold gas in opposition to the flow of external radiation; behind the leading edge of this wave the gas is extremely hot and radiates strongly. The radiation fluxes with the continuous spectrum are of the same order of magnitude as the source flux of radiation, so that a wave of radiation develops.

The propagation of a wave of radiation was considered in [5, 6]. Estimates were given for the parameters of radiation waves propagating in opposition to a flux of external radiation of very great intensity (flux density $\sim 10^5$ MW/cm^2 or over), leading initially to the breakdown of the cold air. The temperature in these waves is $\sim$ 50-100 eV, the velocity $\sim$ 100 km/sec and over. Such radiation waves are supersonic and constitute an alternative to detonation waves [5-7].

The opposite limiting case of the very slow supply of energy to a plasma cloud was considered in [8-12]. For a small diameter of the beam the pressure is able to level itself out and become equal to the surrounding atmospheric pressure. Under such conditions not only processes involving the lateral expansion of the plasma column but also energy losses through the sides of the column play a significant part. In estimating the rate of propagation of the spark allowance is made for the substantial role of molecular and electron heat conduction. The propagation velocities equal several m/sec.

In the plane case here under consideration, the expansion of the heated layer in which the energy of the radiation is released leads to the repulsion of the colder and more transparent gases lying above it at velocities of $\sim$ 1-5 km/sec for flux densities of 1-100 MW/cm^2. A shock wave travels through the gas in opposition to the light beam; the pressure behind it (including the zone of energy release may greatly exceed atmospheric. The amplitude of the shock wave is not too great and the heating of the gas behind the leading edge of the shock wave is not very substantial – the gas remains transparent. The source radiation then penetrates almost without absorption to the edge of the plasma cloud (the front of the radiation wave), which as we shall later show lags with respect to the leading edge of the shock wave for flux densities of the external monochromatic radiation lower than a certain limit, remaining subsonic.

This paper is devoted to an analysis of the propagation of such a subsonic radiation wave, using numerical methods. The solution of transient gasdynamic problems allowing for the transport of continuous-spectrum radiation, traveling in different directions with different quantum energies, involves serious difficulties (mostly of a technical character), even under conditions of one-dimensional geometry. Attempts at the direct numerical integration of the system of equations of radiative gasdynamics, while keeping a fairly strict account of the spectral and angular characteristics of the radiation, are seriously restricted by the limited capabilities of modern electronic computers.

In view of this we shall use the method of averaging the radiation-transport equations with respect to both angles and frequencies. The effectiveness of angular averaging (which was employed earlier in [13-15]) for the case of plane geometry is based on the fact that the average cosine can only vary over a fairly narrow range. There are considerably greater difficulties in allowing for the true spectral composition of the radiation, owing to the complicated frequency dependence of the spectral absorption coefficients. The use of the multigroup approximation [13-15], taking a fairly precise account of the spectral composition, may require too large a number of groups. Here we shall employ averaging over the true spectrum for a small number of groups [16]. Similar ideas were recently expressed in [17]. Since the method of averaging [16] proved to be very effective in relation to the present problem, we shall set out the results of experience accumulated in the use of this method, providing useful guidance for the solution of other problems in radiative gasdynamics.

The system of equations describing the motion of the gas (with both inflow and outflow of energy) in the one-dimensional plane case takes the form

$$\frac{\partial u}{\partial t} + \frac{\partial p}{\partial m} = 0 \tag{1.1}$$

$$\frac{\partial x}{\partial t} = u, \quad \frac{\partial x}{\partial m} = v \tag{1.2}$$

$$\frac{\partial h}{\partial t} - v\frac{\partial p}{\partial t} + \frac{\partial q}{\partial m} = 0 \tag{1.3}$$

$$h = pv\gamma / (\gamma - 1), \quad \gamma = \gamma(h, p), \quad T = T(h, p) \tag{1.4}$$

Here t is the time, x the Euler constant, m the Lagrange mass coordinate, u the velocity, v the specific volume ($v = 1/\rho$, where ρ is the density), p is the pressure, h the specific enthalpy, γ the effective adiabatic index, and T the temperature.

We shall consider that the energy is transferred solely by radiattion, including the radiation of the external source and the "intrinsic" radiation (that generated in the hot gas), so that

$$q = q_l + q_r \tag{1.5}$$

where q_i is the radiation flux density from the source and q_r is the flux density of the intrinsic radiation.

The radiation transfer equation takes the form

$$\mu\frac{\partial I_\varepsilon}{\partial m} = -\varkappa_\varepsilon (I_\varepsilon - B_\varepsilon), \quad \mu = \cos\theta \tag{1.6}$$

Here I_ε is the spectral intensity of the radiation, θ is the angle between the direction of propagation of the radiation and the x axis, $\varkappa_\varepsilon$ is the spectral mass absorption coefficient of the radiatoin, ε is the energy of the quanta, B_ε is the Planck function, defined by the equation

$$B_\varepsilon = \frac{15}{\pi^4}\frac{\sigma\varepsilon^3}{\exp(\varepsilon/T) - 1} \tag{1.7}$$

where σ is the Stefan– Boltzmann constant. The relationships

$$\varkappa_\varepsilon = \varkappa(\varepsilon, h, p) \tag{1.8}$$

characterizing the optical properties of the bases in the range of variation of the parameters under consideration are usually specified in tabulated form.

We shall consider that the radiation of the source is directed along the x axis, for the sake of determinacy in the negative sense

$$\frac{\partial q_l}{\partial m} = \varkappa_l q_l \tag{1.9}$$

Here $\varkappa_l$ is the mass-absorption coefficient of the source radiation $\varkappa_l = \varkappa_\varepsilon(\varepsilon_l)$, where ε_l is the energy of the source quanta.

Subsequently we shall use the characteristics of the continuous-spectrum radiation integrated with respect to both angles and spectral energies

$$U_\varepsilon^\pm = \int_{\mu_1^\pm}^{\mu_2^\pm} I_\varepsilon d\mu, \quad q_\varepsilon^\pm = 2\int_{\mu_1^\pm}^{\mu_2^\pm} I_\varepsilon \mu d\mu \tag{1.10}$$

$$\mu_1^+ = 0, \quad \mu_2^+ = 1, \quad \mu_1^- = -1, \quad \mu_2^- = 0$$

$$U_i^\pm = \int_{\varepsilon_1^i}^{\varepsilon_2^i} U_\varepsilon^\pm d\varepsilon, \quad q_i^\pm = \int_{\varepsilon_1^i}^{\varepsilon_2^i} q_\varepsilon^\pm d\varepsilon \tag{1.11}$$

For the sake of convenience we shall in future omit the indices plus and minus attached to the limiting cosines $\mu_1\pm$ and $\mu_2\pm$ and the index i on the group limits $\varepsilon_1{}^i$ and $\varepsilon_2{}^i$.

We have

$$q_r = \sum_i (q_i^+ + q_i^-) \tag{1.12}$$

2. The continuous-spectrum radiation may be characterized at every point by a spectral-direction diagram $\psi_\varepsilon\pm$ and a spectral average cosine $c_\varepsilon\pm$.

$$\psi_\varepsilon\pm = I_\varepsilon/U_\varepsilon\pm, \quad c_\varepsilon\pm = \int_{\mu_1}^{\mu_2} \mu\psi_\varepsilon\pm d\mu \tag{2.1}$$

and also by a radiation spectrum $(\varphi_\varepsilon\pm)_i$ in each group, a unilateral average group cosine $c_i\pm$, and a unilateral average group absorption coefficient $\langle\varkappa\rangle_i\pm$

$$(\varphi_\varepsilon)_i\pm = U_\varepsilon/U_i\pm \tag{2.2}$$

$$c_i\pm = \int_{\varepsilon_1}^{\varepsilon_2} c_\varepsilon\pm (\varphi_\varepsilon)_i\pm d\varepsilon \tag{2.3}$$

$$\langle\varkappa\rangle_i\pm = \int_{\varepsilon_1}^{\varepsilon_2} \varkappa_\varepsilon (\varphi_\varepsilon)_i\pm d\varepsilon \tag{2.4}$$

When the radiation spectrum approaches the Planck form

$$(\varphi_\varepsilon)_i{}^p = B_\varepsilon/B_i, \quad B_i = \int_{\varepsilon_1}^{\varepsilon_2} B_\varepsilon d\varepsilon \tag{2.5}$$

the absorption coefficient $\langle\varkappa\rangle_i\pm$, averaged over the true spectrum approaches the average Planck group absorption coefficient $B_i(T)$ defined thus

$$\varkappa_i{}^p = \int_{\varepsilon_1}^{\varepsilon_2} \varkappa_\varepsilon (\varphi_\varepsilon)_i{}^p \, d\varepsilon \tag{2.6}$$

Tables of the Planck group absorption coefficient $\varkappa_i{}^p$ and the function $B_i(T)$ may be set up before solving the problem.

Having integrated the transport equation (1.6) with respect to μ and ε we obtain [16]

$$\frac{\partial q_i\pm}{\partial m} = -\frac{\xi_i\pm}{c_i\pm}(\varkappa_i^\circ)\pm q_i\pm + 2\varkappa_i{}^p B_i \tag{2.7}$$

$$\xi_i\pm = \langle\varkappa\rangle_i\pm / \varkappa_i^\circ, \quad (\varkappa_i^\circ)\pm = \varkappa_i^\circ(h, p) \tag{2.8}$$

where $\xi_i\pm$ is the "distortion" coefficient and $\varkappa_i{}^\circ$ the "reference" absorption coefficient, averaged, for example, over the Planck or some other standard spectrum. For the sake of convenience, we omit the plus and minus indices on $\varkappa_i{}^\circ$, on the right-hand sides of (2.8) and subsequently. We note that the definition of $\varkappa_i{}^\circ$ may differ in different regions of the problem.

The direct solution of transient problems using the complete system of equations of radiative gas dynamics, i.e., integrating the transport equation (1.6) at each time step, is practically impossible if a sufficient number of ε and μ points are to be taken.

However, in many cases there is no need to attempt this. Usually the average characteristics of the field of radiation change more slowly than the directional diagram and the radiation spectrum at individual μ and ε points. The gas-dynamic problem may therefore be solved by using averaged transport equations, for example, in the form (2.7) and (2.8), the solution of (1.6) being carried out more rarely – at certain specific instants of time (the instants of "averaging") – with the aim of determining the spectral and angular characteristics of the radiation, if these are of independent interest, and finding the coefficients $\xi_i\pm$, $c_i\pm$ in the averaged equations (2.7).

The difficulties arising in a specific realization of the averaging method are due to the fact that the average coefficients are functions of two variables (m and t), while at the instant of averaging they are functions of one variable. It was proposed to determine the time dependence by a recalculation process in [16].

The $\xi_i\pm$ relationship, for example, may be expressed in the form $\xi_i\pm(\omega_i\pm, t)$, where $\omega_i\pm(m, t)$ is a certain new variable chosen in such a way as to make the time dependence (dependence on t) as weak as possible. If we make a successful choice of $\varkappa_i°$ and the "principal variable" $\omega_i\pm$, the time variation of $\xi_i\pm$ for $\omega_i\pm$ = const will be insignificant, and the averaging may be carried out quite infrequently. The choice of $\varkappa_i°$ and $\omega_i\pm$ is not unique. A final choice has to be made on the basis of a preliminary analysis of the specific problem, and may be corrected during the actual calculations.

This is also true of the choice of the number of groups, Too great an increase in the number of groups increases the amount of information to be stored and the volume of calculations based on the averaged transport equations, ultimately leading to the necessity of solving the spectral problem in every time layer. The number of groups must nevertheless not be so low as to be incapable of allowing fpr the specific characteristics of the problem and depicting its qualitative aspects correctly.

3. Let us consider the expansion of the gas in which the energy of the source radiation is released, and also its redistribution by virtue of the transfer of radiation from the continuous spectrum. Let the velocity of sound in the hot gas be high and let the expansion of the boundary of the plasma cloud take place fairly slowly. During the period of energy inflow the sound perturbations will therefore be able to travel repeatedly through the hot volume, equalizing the pressure in the latter, i.e., in the hot region we may put

$$\frac{\partial p}{\partial m} = 0, \quad p = p^{\circ}(t) \tag{3.1}$$

This enables us to avoid accounting for the high-frequency pressure perturbations (acoustic vibrations), which, as direct numerical calculations show, travel rapidly through the heated layer, without changing the average pressure and other characteristics of the radiation wave. The amplitude and position of these perturbations are to some extent random and there is no need to take an exact account of them.

Subsequently we shall solve the problem subject to the assumption (3.1). Let us now set out the method of solving the problem.

Let the pressure $p°(t)$ be given. Then the system in which Eqs. (1.1) and (1.2) are discarded enables us to solve the problem of radiative energy transfer and find h(m, t) and v(m, t). From Eqs. (1.2) we find the Euler coordinate x(m, t) and the velocity u(m, t) of all the particles, including the boundaries of the hot volume, i.e., the velocity of the "piston" $u°(t)$. By solving the comparatively simple problem of the motion of the cold layers under the action of such a piston, allowing for a possible fall in pressure and the motion of the shock wave, i.e., by using (1.1) but not taking account of the energy evolution in these transparent layers, we may find the values of the pressure everywhere, even in the piston, and hence in the hot region. The solutions of the external and internal problems have be be matched ot one another.

When the external flux of radiation and the pressure are constant in time, the gas velocity u_0 at the boundary of the absorbing layer (the leading edge of the radiation wave) also remains constant. Then the velocity of that part of the gas which is cold and transparent to the source radiation remains constant up to the leading edge of the shock wave, which moves at a constant velocity a and frontal pressure p_s ($u° = u_s$ and $p° = p_s$). If the shock wave is strong, we obtain

$$p^{\circ} = A(\gamma, \gamma_s)\, q_0^{2/3}\, \delta^{1/3}, \quad u^{\circ} = B(\gamma, \gamma_s)\, q_0^{1/3}\, \delta^{-1/3} \tag{3.2}$$

where γ is the adiabatic index in the hot layer, γ_s is that in the cold layer behind the leading edge of the shock wave, δ is the density of the gas in front of the leading edge, referred to the normal density ρ_0 of air, q_0 is determined by the equation

$$q_0 = q_l^{\circ} + q_r^{\circ} \tag{3.3}$$

Here $q_l°$ is the flux density of the external radiation incident upon the leading edge of the radiation wave, $q_r°$ is the flux density of the intrinsic radiation (continuous spectrum) moving away from the leading edge.

If p is in bar, u in km/sec, q in MW/cm^2, $\gamma = 1.18$ and $\gamma_s = 1.4$, then $A \approx 13$ and $B \approx 1$.

Thus to an accuracy disregarding the radiation losses $q_r°$ (which are not known in advance) we may use Eq. (3.2) to estimate the pressure in the layer without solving the transport problem. As we shall later show, $q_r°$ is small compared with $q_l°$.

The temperature and radiation fluxes in the radiation wave and the mass velocity of the propagation of its leading edge depend only slightly on the pressure. Hence subsequent corrections make little difference to the initial estimte of $p°$.

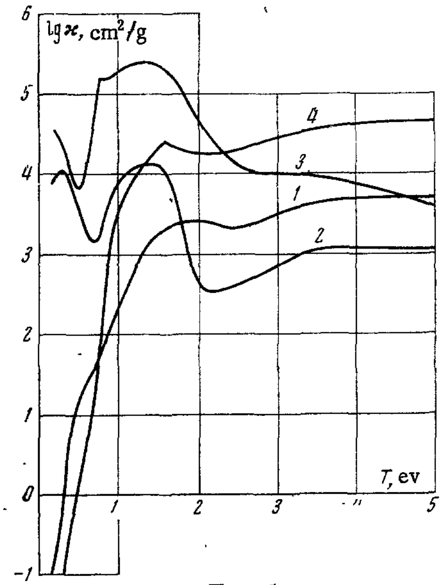

Fig. 1

In accordance with the general presentation of the problem, the propagation of the radiation wave is calculated on the assumption that a plasma layer opaque to incident radiation has already formed at t = 0. The time required for the formation of this layer, measured from the initial action of the source, is determined by the theory of [2-4]. In solving (1.6) and (2.7) we assume that no continuous-spectrum radiation falls on the boundaries of the volume under consideration from outside.

4. By way of example let us give the results of some numerical calculations relaing to a subsonic radiation wave traveling under the influence of monochromatic radiation with a quantum energy of ε_l = 1.16 eV in air of normal density. For this case ($\rho_0 = 1.29 \cdot 10^{-3}$ g/cm^3) with $q_l{}^\circ$ = 5, 10 and 20 MW/cm^2 respectively, according to Eq. (3.2) we obtain p° = 30, 50, and 80 bar (allowing for the radiative energy losses determined in the calculation).

The thermodynamic properties of the air were determined from tables [18].

In calculating the transfer of radiation in hot air we used the tables of [19]. For temperatures of over $20 \cdot 10^3$ °K and quantum energies $\varepsilon > 18.5$ eV (at all temperatures) these tables were supplemented by the results of calculations based on the method of [19], but without allowing for the lines. In order to reduce the amount of calculations these tables were transformed by uniting mutually adjacent spectral intervals. The number of spectral intervals was thus reduced to 175.

The intensity of the radiation was determined for 24 rays.

Averaging was carried out in three groups. The first group incorporated quanta with energies of $\varepsilon < 6.51$ eV (long-wave radiation). The third group contained radiation with $\varepsilon > 7.75$ eV. The second group was intermediate.

Figure 1 illustrates the temperature dependence of the average Planck absorption coefficients and also the monochromatic-radiation absorption coefficient for a pressure of p° = 50 bar. The figures 1-4 denote the curves of $\varkappa_{1}p$, $\varkappa_{2}p$, $\varkappa_{3}p$, and $\varkappa_l$ remain large; the quanta of these groups cannot travel to very great distancse from the hot zone. These provide for the heating of the cold layers and the forward motion of the radiation wave in the opposite direction to the radiation flux of the source. The qualitatively different type of behavior of $\varkappa_\varepsilon(T)$ for soft and hard quanta as $T \to 0$ necessitates the introduction of at least two groups.

Actually the boundary separating those parts of the spectrum with relatively long and relatively short ranges changes with temperature [19]. This may best be allowed for by introducing an intermediate group. By comparing $\varkappa_{2}p$ and $\varkappa_{3}p$ with $\varkappa_l$ we may convice ourselves of the fact that the region with low temperatures is transparent to the external radiation, but starting from a temperature of 1-2 eV this radiation is intensively absorbed in the radiation wave.

For all temperatures $\varkappa_{2}p < \varkappa_{3}p$. The quanta of the intermediate group thus form heating "temperature tongues" in front of the leading edge of the radiation wave. Such tongues may clearly be seen in Fig. 2, which gives the temperature distributions at various instants of time (up to 7 μsec) for the case in which the flux density $q_l{}^\circ$ = 10 MW/cm^2 (the following figures also relate to this case). The numbers 0-4 correspond to the instants of time 0, 2.8, 4.3, 5.8, 7.3 μsec. We see from Fig. 2 that the leading edge of the radiation wave travels toward the flow of radiation from the source at an almost constant velocity (the mass flow through the leading edge is m_r = 25 g/cm^2/sec). In addition to this, a trailing edge develops, moving in the opposite direction at a lower velocity and gradually becoming slower. The maximum temperature ~ 3.5 eV varies slowly with time.

Figure 3 shows the distribution of the unilateral group flux densities of the radiation $q_i \pm (i = 1, 2, 3)$ with respect to the same coordinate at the instant of time t = 2.8 μsec. The numbers 1-3 denote the relationships for $q_1{}^+$, $q_2{}^+$, $q_3{}^+$, I-III those for $q_1{}^-$, $q_2{}^-$, $q_3{}^-$. We see that the quanta of the first ("soft") group pass out of the hot zone and cease being absorbed. The quanta of the third ("hard") group generated in the hot air are almost completely absorbed in the narrow layers constituting the "fronts" (leading and trailing edges) of the radiation wave.

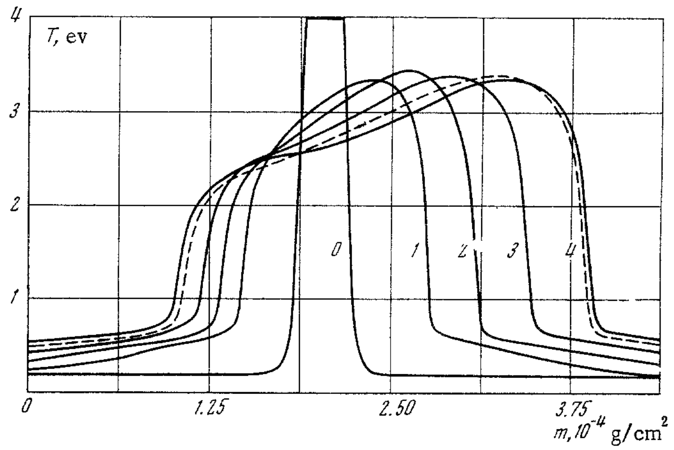

Fig. 2

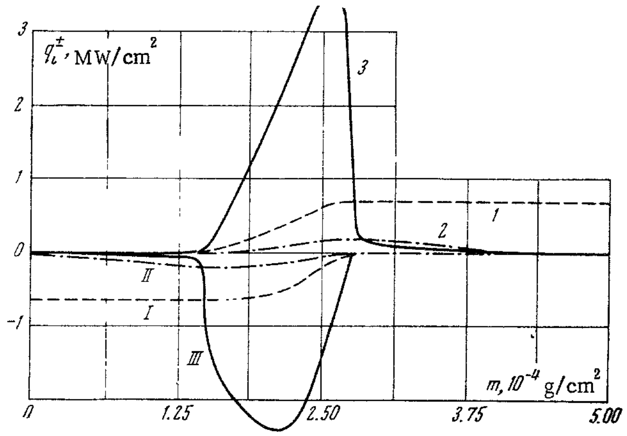

Fig. 3

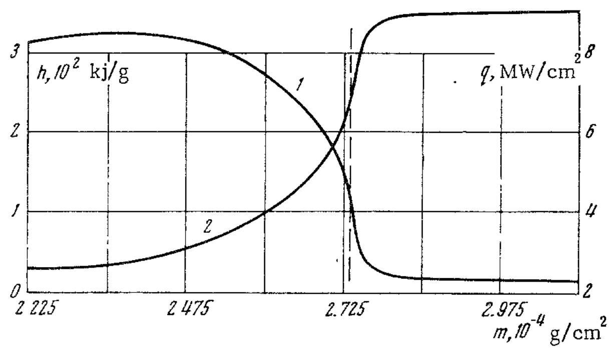

Fig. 4

The flux density of the intrinsic radiation emerging from the hot zone through the leading and trailing edges of the radiation wave is relatively low, 1 MV/cm^2, or only 10% of $q_l°$.

Figure 4 presents a detailed distribution of the parameters h (curve 1) and -q (curve 2) with respect to mass close to the leading edge of the radiation wave at the instant t = 2.8 μ sec. The nominal position of the leading edge, defined as the point at which the absolute value of $\partial h/\partial m$ reaches a maximum, is indicated by the broken line. The distribution of the computed points with respect to mass was nonuniform; for a total number of such points equal to 240 (in the present version) some 40 points with a mass interval five times smaller than in the hot region were placed at the leading and trailing edges.

The total number of computed layers in the problem was quite large (up to the instant of time 7 μsec)—about 700.

Over this time four averaging were carried out (at the instants 1.8, 2.8, 4.8, and 7.3 μsec). In order to illustrate the influence of the recalculations procedure the results obtained in the preliminary stage of calculation are shown as a broken line in Fig. 2 for the latter instant of time. Satisfactory agreement with the results obtained after recalculation (continuous curves) was also obtained for the other instants of time, convincing us that the averaging frequency adopted in the present analysis was sufficient. The long intervals between averagings demonstrate the effectiveness of our use of the averaging method in connection with the present problem.

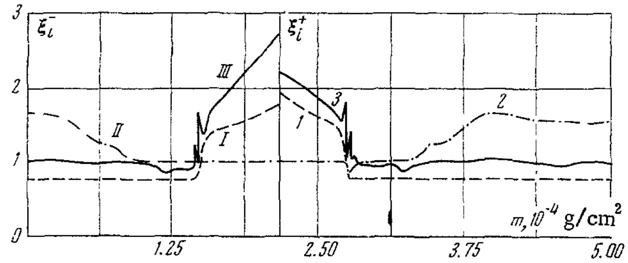

Fig. 5

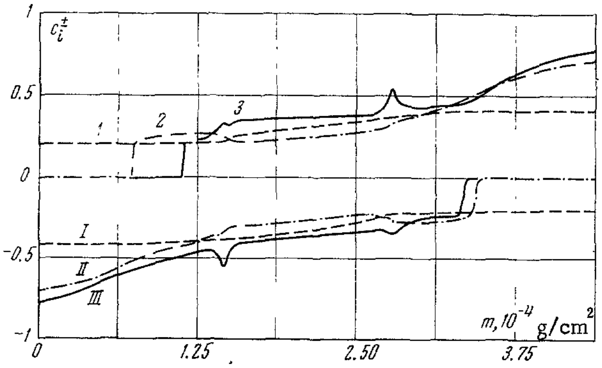

Fig. 6

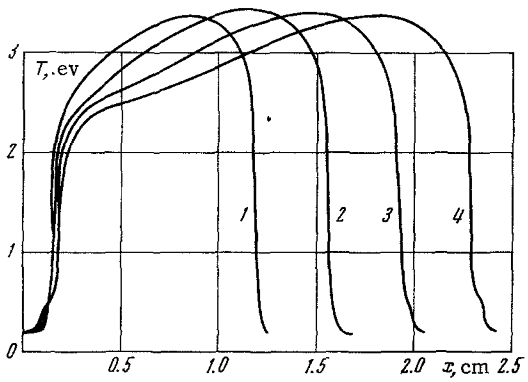

Fig. 7

Let us give some more detailed attention to the choice of reference coefficients. In the present problem there are three regions with different properties: a hot zone behind the leading edge of the radiation wave with a temperature close to maximum, and two regions of cold air in front of the corresponding edges of the radiation wave ("tongues"). In the first region the reference coefficient adopted was the average Planck absorption coefficient. In front of the leading and trailing edges of the radiation wave the spectral composition of the radiation varies little by comparison with the spectrum at a "frontal point" but differs considerably from the Planck spectrum for the temperature existing in front of the edge. Hence the value of $\varkappa_i^p$ in the cold region is considerably smaller than the true value of $\langle\varkappa_i\rangle$ (for the first group by several orders of magnitude). (Correspondingly ξ_1 would be of the order of 10^3–10^4.)

The sharp change in ξ_i which occurs on passing through the leading edge of the radiation wave makes interpolation difficult and reduces the accuracy of numerical calculations. In the problem under consideration we may relate the reference absorption coefficient in the tongues to the spectrum at a certain depth i inside the leading edge of the radiation wave. On passing through the leading edge there is a change in the reference absorption coefficient, and the value of ξ_i varies less rapidly. This may be seen from Fig. 5 which gives the mass dependence of $\xi_i\pm$ at the instant t = 2.8 μsec. Figure 6 shows the $c_i\pm(m)$ relationship for the same instant of time. In the last two figures the serial number of i of the corresponding group is denoted by an arabic figure for the relationships with index plus and a Roman figure for those with index minus. We see that the range of variation of c_i is small.

By way of "principal" variable $\omega_i\pm$ we chose the group optical thickness, determined by the equations

$$d\tau_i^\pm = \pm\varkappa^\circ_i dm$$

and reckoned from the leading and trailing edges of the radiation wave.

Figure 7 shows the temperature distribution T and Fig. 8 shows the flux density distribution of the incident radiation q_l at different instants of time t in Euler coordinates x. Curves 1–4 correspond to the

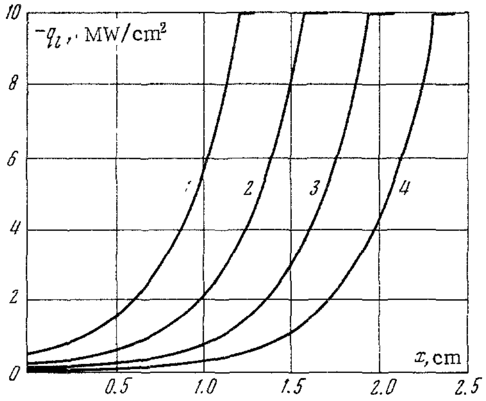

Fig. 8

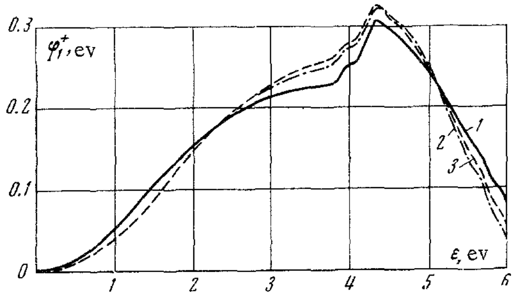

Fig. 9

instants of time 2.8, 4.3, 5.8, 7.3 μsec. In these coordinates the edges are sharper than in Lagrange coordinates (Fig. 2), since as the temperature falls the density increases. The rate of propagation of the leading edges D ≈ 2.5 km/sec. At the same time the hot zone, in which the density is low, is drawn out to relatively long distances. The characteristic depth of penetration of the radiation (the distance in which the quantity q_i falls by a factor of e as compared with $q_l°$) is 0.3 cm. The flux of radiation q_l with an amplitude 10-20% of the value of $q_l°$ also exerts an influence to a considerable depth – layers of the order of 1.0-1.5 cm thick. The point of maximum temperature lies at about 0.3 cm from the leading edge; the temperature behind this varies very little throughout almost the while thickness of the radiation wave ("plateau").

Figure 9 shows the spectrum of the radiation emitted from the leading edge of the radiation wave in air at various instants of time t = 2.8, 4.8, and 7.3 μsec (curves 1-3 respectively); it is quite smooth and approximately corresponds to the Planck distribution for a temperature of ~1.5 eV. Inside the radiation wave the spectra are more complicated, especially in the second group and close to the lower limit of the third group.

The picture of radiation-wave propagation for slightly greater and smaller flows is similar to the foregoing for the case $q_l° = 10$ MW/cm². The maximum temperature depends only slightly on the flux density; for $q_l° = 5$ and 20 MW/cm² it is respectively 3 and 4 eV. After the period of establishment the leading edge of the radiation wave moves at an almost constant velocity. This velocity varies approximately as indicated by Eq. (3.2).

The mass flow m_r through the leading edge of the radiation wave is much smaller in the range of flux densities considered than that through the edge of the shock wave propagation in front of the radiation wave. Thus for $q_l° = 20$ MW/cm² we have $m_r = 60$ g/cm²/sec, whereas $m_s = 340$ g/cm²/sec. Thus the radiation wave lags on the shock wave.

The temperature at the leading edge of the shock wave is 3000° K; according to [19] such a shock wave is transparent for very great thicknesses of the shock-compressed layer.

For a temperature of $\sim 35\text{-}45 \cdot 10^3$ °K (inside the radiation wave) the velocity of sound is ~10 km/sec, which considerably exceeds the linear rate of propagation of the leading edge of the radiation wave and justifies the foregoing (§ 3) assumption as to the uniformity of the pressure in the hot zone.

For flux densities greater than those considered, the process here described is replaced by one of ligh detonation [5-7] when the absorption zone moves at the same rate as the shock-wave front. For very high densities the flux of the radiation wave over takes the shock wave, becoming supersonic [5, 6].

The latter case is qualitatively reminiscent of the replacement of a temperature wave of the second kind (TW-II) by a temperature wave of the first kind (TW-I) [20].

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