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CHANGE IN AMPLIFICATION FACTOR IN THE SHOCK LAYER WHEN A SUPERSONIC FLOW WITH AN INVERTED POPULATION FLOWS AROUND BLUNT BODIES

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INTRODUCTION

In modeling the supersonic flow of a relaxing gas around a solid body it is important to make a detailed physicochemical analysis of the internal structure of the flow. The working gas used to simulate the real flow often comprises gas mixtures obtained by the combustion of hydrocarbon fuels and contains CO_2, N_2, O_2, and H_2O molecules. Of special interest in simulation problems is the circumfluence of a nonequilibrium flow with an inverted population of the vibrational levels of the CO_2 molecules. A calculation of the amplification factor for the (00⁰1) – (10⁰0) transition of the CO_2 molecule during the development of indirect jumps in compression (shock waves) in an inverted medium was presented in [1]; there was a reduction in amplification factor for the vibrational–rotational transition P(20) over the pressure range in which the greatest contribution to spectral-line broadening was due to the collision mechanism. A fall or rise in amplification factor was observed in [2], according to the intensity of the shock wave and the rotational quantum number.

In this paper we shall study the changes taking place in the amplification factor when blunt solids are immersed in a gas flow (both in the subsonic and in the supersonic parts of the shock wave) as the angle of inclination of the shock wave to the direction of the incident flow varies from 90° to the Mach angle; we shall also study the influence of small perturbations traveling through the inverted medium on the amplification factor.

§1. It is well known [3] that for the vibrational–rotational transition (00⁰1) – (10⁰0) the amplification factor of a weak signal may be written

$$G = (\lambda^2 A_{nm}/8\pi\sqrt{\pi}c)[N_n - (g_n/g_m)N_m](a/\Delta_c)H(a, 0), \tag{1.1}$$

where λ is the wavelength of the transition; A_{nm} is the Einstein coefficient for the spontaneous transition $n \to m$; c is the velocity of light; the parameter $a = (\Delta_c/\Delta_D)\sqrt{\ln 2}$; Δc is the half-width of the line accounted for by collisions; Δ_D is the Doppler half-width; N_n, N_m, g_n, g_m are the populations and statistical weights of the upper and lower levels, respectively; and $H(a, 0)$ is the Voigt function in the center of the line. The temperature dependence of the shock half-width was taken as proportional to $T^{-1/2}$.

Let us consider the axisymmetrical passage (around a cylinder with spherically blunted ends) of a nonviscous supersonic homogeneous flow of relaxing gas mixture with an inverted population in the incident flow

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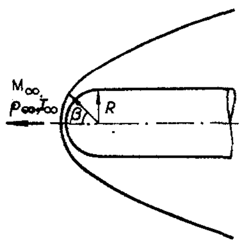

Fig. 1

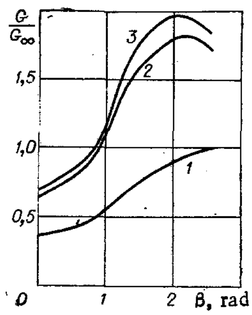

Fig. 2

between (00^01) and (10^00) levels of the CO_2 molecule. The principal shock wave is formed in front of the solid. In the shock wave there is a sudden rise (extending over several mean free paths of the molecules) in density and in the temperature of the active degrees of freedom of the molecules comprising the gas mixture. The internal degrees of freedom, however, have a considerably greater relaxation length. Thus, on passing through the leading edge of the shock wave the populations of the vibrational levels are preserved intact.

In the frontal part of the body the shock wave is almost at right angles to the direction of the incident flow; then it curves and ultimately passes into the Mach line.

The amplification factor of a small signal in the shock layer was calculated by the method of streamlines. The essence of this method lies in integrating the equations of gasdynamics and the equations of vibrational relaxation of the gas mixture in question along the streamlines for a known pressure distribution with respect to the latter and a specified shape of the principal shock wave. In the calculations we used the coordinate system S, ψ (ψ is the stream function; S is the distance along the streamline). The origin $S = 0$ was taken at the jump of compaction. Corresponding to each streamline is a specific angle β characterizing the point of transition through the shock wave (Fig. 1). The shock wave is regarded as infinitely thin, so that in passing through it the composition of the gas remains unchanged. The position and shape of the shock wave are taken from [4, 5].

The geometry of the streamlines and the pressure distribution along them appear in [6]. The pressure, density, and enthalpy (velocity) in the shock wave are completely determined in accordance with the Rankine–Hugoniot conditions [7].

The parametric calculations so executed enable us to explain the behavior of the amplification factor at the leading edge of the shock layer for various mechanisms of spectral-line broadening. Figure 2 gives the results of some calculations regarding the changes taking place in the amplification factor of a small signal for the P(20) transition at the leading edge of a shock wave formed by the flow of a $CO_2-N_2-H_2O$ gas mixture around a thin cylinder with spherical ends (radii of spheres $R = 1.5$ cm). The horizontal axis gives the values of the angle β (Fig. 1). The calculations were carried out for a mixture of the following composition: $\alpha_{CO_2} = 0.1$, $\alpha_{N_2} = 0.89$, $\alpha_{H_2O} = 0.01$ (concentrations given in molecular parts). The conditions in the incident flow were taken from the calculations of [1, 8].

The Mach number of the incident flow $M_\infty = 6$; the temperature $T_\infty = 200°K$. Curve 1 corresponds to an incident flow density of $\rho_\infty = 10^{-4}$, 2 to $\rho_\infty = 10^{-6}$, 3 to $\rho_\infty = 10^{-7}$ g/cm^3. (The corresponding values of a_∞ lie in the following ranges: $a_\infty > 1.4$; $0.2 < a_\infty < 1.4$; $a_\infty < 0.2$.

In the high-density region in which the spectral-line broadening is due to collisions both in the incident flow and at the leading edge of the shock wave, the parameter $a \gg 1$, and the ratio of the amplification factor behind the jump to its value in front of the jump (G/G_∞) is independent of density, since the rise in G/G_∞ due to the increase in the number of active particles in the shock layer is completely compensated by the reduction in G/G_∞ due to collision broadening. The change in G/G_∞ is here largely influenced by the translational temperature associated with the populations of the rotational levels of the CO_2 molecules (curve 1, Fig. 2). Since the shock wave formed in front of the body has different angles of inclination σ relative to the incident flow, $\arcsin(1/M_\infty) < \sigma \leq \pi/2$, the increments in the translational temperature at the leading edge of the shock wave will also be different. In the shock wave adjacent to the frontal part of the body the temperature rise is at a maximum; a long way from the body it is at a minimum. The value of the quantum number $I_{max} = \sqrt{kT/hcB}) - 1/2$, corresponding to the maximum in the distribution over the rotational levels, increase with rising temperature. Hence for fixed large values of the rotational quantum number I_n the temperature rise at the jump may lead to an increase in the inversion of populations as I_{max} approaches I_n.

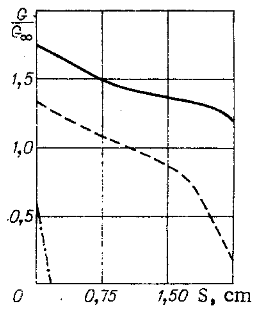

Fig. 3

On further increasing the translational temperature the degree of inversion starts falling. However, the expansion of the spectral line increases monotonically with rising temperature. If the contribution from the inversion behind the jump to the amplification factor exceeds the contribution due to the spectral-line broadening, we have $G/G_\infty > 1$.

It may be shown that on satisfying the condition

$$\alpha = (\Theta/T_\infty)I_n(1 + I_n) > 1, \text{ where } \Theta = hcB/k, \tag{1.2}$$

the curve $G/G_\infty = f(\sigma)$ has a maximum; the inclination of the shock wave for which the maximum change in G/G_∞ is realized at the front of the jump (for fixed conditions in the incident flow) may be obtained from

$$\sigma = \arcsin\left\{\frac{(\varkappa-1)^2 - 4\varkappa + (\varkappa+1)^2\alpha + \sqrt{[(\varkappa-1)^2 - 4\varkappa + (\varkappa-1)^2\alpha]^2 + 16\varkappa(\varkappa-1)^2}}{4\varkappa(\varkappa-1)}\right\}^{1/2},$$

where $\varkappa$ is the effective adiabatic index. If condition (1.2) is not satisfied, the change in G/G_∞ along the shock wave will bear a monotonic character.

We see from Fig. 2 (curve 1) that for small values of the angle β the ratio $G/G_\infty < 1$, since as a result of the intensive rise in temperature close to the frontal surface of the solid a sharp broadening of the spectral line takes place. With increasing angle β, i.e., falling intensity of the shock wave, the role of broadening diminishes and there is a rise in the contribution to the amplification factor arising from the increment in the population of the rotational level $I_n = 19$. The effect is that G/G_∞ tends toward unity.

For low densities in the incident flow (when $a \ll 1$, i.e., a Doppler line contour is realized) an increment in G/G_∞ at the leading edge of the shock wave takes place by virtue of an increase in the population of the rotational levels, and a reduction in G/G_∞ as a result of Doppler broadening.

The maximum temperature increment takes place in the part of the shock wave lying nearest to the critical streamline, with a corresponding change in line broadening, which is proportional to $\sqrt{T}$. The increase in density cannot compensate the reduction in G/G_∞ due to the line broadening. Hence in this region $G/G_\infty < 1$. For smaller inclinations of the shock wave the jump in translational temperatures at the leading edge diminishes; this is responsible for the smaller reduction in G/G_∞ on account of the line broadening. On further reducing the intensity of the shock wave, the contribution to G/G_∞ due to the increase in the population of the rotational level starts exceeding the reduction in G/G_∞ due to the Doppler broadening of the line, and G/G_∞ rises above unity. For a low intensity of the shock wave ($\beta > 1$) on account of the slower rise in density the increase in the population of the level $I_n = 19$ cannot any longer compensate the influence of the line-broadening factor, and the ratio G/G_∞ starts falling (curve 3, Fig. 2). If the coefficient a lies in the range $1.4 > a > 0.2$ both in the incident flow and in the shock layer, the broadening of the spectral line is affected by both the collisions and the Doppler effect. In this case there may be either a rise or a fall in the amplification factor at the leading edge of the shock wave, depending on the inclination of the shock wave to the velocity vector of the incident flow (curve 2, Fig. 2).

§2. The increase in density and translational temperature behind the leading edge of the shock wave leads to a change in the character of the relaxation processes in the shock layer, in which (as a result of the equalization of the translational and vibrational temperatures) slight changes in temperature and density and substantial changes in the population of the levels take place, leading to a further change in the amplification factor.

In calculating the populations of the vibrational levels of the molecules in the shock layer we used a system of kinetic equations allowing for processes of multiple-quantum vibrational–vibrational exchange (with

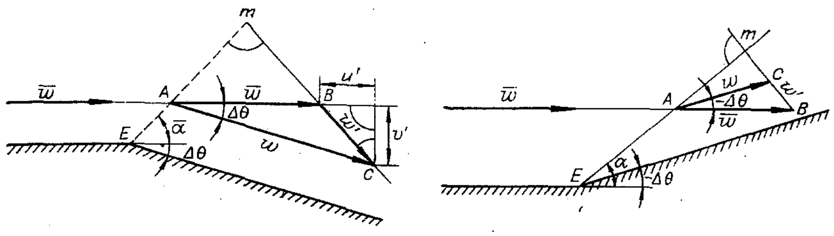

Fig. 4

participation of degenerate modes) and for the contribution of transitions between highly excited vibrational states [8]. The choice of such a system of kinetic equations arose from the fact that any calculation of the relaxation zone behind the shock wave based on the approximate kinetic equations leads to serious errors in the distribution of vibrational temperatures and the populations of the vibrational levels of the CO_2 molecules [8].

Figure 3 shows the variations in amplification factor along the streamlines in the relaxation zone of the shock layer for the following conditions in the incident flow: $M_\infty = 6$, $\rho_\infty = 10^{-6}$ g/cm^3, $T_\infty = 200$°K, $G_\infty = 0.475 \cdot 10^{-3}$ cm^{-3}. Here S is the length of the streamline; S = 0 corresponds to the shock wave. The dashed-dot line corresponds to the streamline passing through the direct shock wave ($\beta = 0°$) (Fig. 1), the dashed curve to the streamline defined by the angle $\beta = 70°$, and the continuous curve to the streamline defined by the angle $\beta = 110°$.

The calculations show that as a result of the equalization of temperatures, there is always a fall in the amplification factor for the vibrational–rotational transitions $(00^01)-(10^00)$ in the relaxation zone behind the principal shock wave. However, in contrast to the flows behind the direct shock wave [2], the substantial fall in temperature along the principal shock wave and also the large tangential component of the velocity of the incident flow have the effect that the gas entering into the shock layer in the supersonic region of circumfluence retains the inverted population for fairly large distances along the streamline. This leads to the development of a region of high amplification factor in the supersonic region of the shock layer (continuous curve in Fig. 3). In the subsonic region of the shock layer there is practically always a fall in the amplification factor, both at the leading edge of the shock wave and in the actual shock layer (dashed-dot curve in Fig. 3). Thus if the conditions of the medium with inverted population in the incident flow promote a rise in the amplification factor of a weak signal at the leading edge of the shock wave, then for low intensities of the shock wave there will be a fairly wide zone with an increased amplification factor in the shock layer.

§3. Let us consider the change in the amplification factor of a weak signal for the $(00^01)-(10^00)$ transition of the CO_2 molecule when a weak shock wave or a wave of rarefaction acts upon the flow. Let us consider the flow around an obtuse angle close to the vertex. In this case the direction of the main flow parallel to the x axis is deviated by a small angle $\Delta\theta$. The angle $\Delta\theta > 0$ for a convex angle and $\Delta\theta < 0$ for a concave angle (Fig. 4: above, rarefaction; below, compression). Here $\bar{w}$ and w are the absolute values of the flow velocities before and after perturbation; w' is the change in the velocity of the main flow.

Since the change in the states of the vibrational degrees of freedom of the molecules takes place on a far larger scale than that of the gasdynamic parameters during the perturbation of the flow, the flow close to the point 0 may be regarded as frozen and the Mach line m as the Mach line of the frozen flow. Using the results of [7] as regards linearized flow around an obtuse angle, for the additional downstream pressure on the line m we obtain

$$p' = -\rho_\infty \frac{\bar{w}^2\Delta\theta}{\sqrt{M_\infty^2 - 1}},$$

where $M_\infty = \bar{w}/\sqrt{\varkappa (k/m) T_\infty}$ is the frozen Mach number. The increments in the progressive temperature and density for the CO_2, N_2, H_2O mixture have the form

$$T' = -\frac{\bar{w}^2\Delta\theta}{\sqrt{M_\infty^2 - 1}}\left[\alpha_{CO_2}\frac{7k}{2m} + \alpha_{H_2O}4\frac{k}{m} + \alpha_{N_2}\frac{7}{2}\frac{k}{m}\right]^{-1}; \qquad (3.1)$$

$$\rho' = -\frac{\rho_\infty^2\bar{w}^2\Delta\theta}{\frac{k}{m}T_\infty\sqrt{M_\infty^2 - 1}}\left(\frac{5 + \alpha_{H_2O}}{7 + \alpha_{H_2O}}\right).$$

In the case of a purely Doppler mechanism of spectral-line broadening, for the P branch of the vibrational–rotational transitions and for $\Delta_c \sim T^{-1/2}$ Eq. (1.1) may be reduced to the following form after linearization of the gasdynamic quantities with respect to the parameter $\Delta\theta$:

$$\frac{G}{G_\infty}=1-\frac{1}{2}\frac{T'}{T_\infty}+\left[\frac{\Theta I_n(1+I_n)}{T_\infty}-1\right]\frac{T'}{T_\infty}+\frac{\rho'}{\rho_\infty}-\frac{\eta}{\left[\exp\left(-\frac{\Theta_3}{T_3}+\frac{\Theta_1}{T_1}+\eta\right)-1\right]}\frac{T'}{T_\infty}, \tag{3.2}$$

where $\eta=2(1+I_n)\Theta/T_\infty$; T_3 and T_1 are the vibrational temperatures of the asymmetrical and symmetrical vibrations of the CO_2 molecules; Θ_3 and Θ_1 are their respective characteristic temperatures.

Since the existence of an inverted population $\Theta_1/T_1-\Theta_3/T_3>0$ is assumed in the unperturbed flow, while for the translational temperatures realized in the nozzles of aerodynamic installations $\eta\ll 1$, the expression $\frac{\eta}{\left[\exp\left(-\frac{\Theta_3}{T_3}+\frac{\Theta_1}{T_1}+\eta\right)-1\right]}\ll 1$, and the last term in Eq. (3.2) may be omitted.

The second term on the right-hand side of (3.2) describes the change in G/G_∞ accounted for by the broadening of the spectral line; the third and fourth describe the change in population due to the change in temperature and density.

On using Eq. (3.1), Eq. (3.2) may be reduced to the form

$$\frac{G}{G_\infty}=1-\frac{M_\infty^2\Delta\theta}{\sqrt{M_\infty^2-1}\,(5+\alpha_{H_2O})}\left[2\frac{\Theta}{T_\infty}I_n(1+I_n)+2+\alpha_{H_2O}\right]. \tag{3.3}$$

It follows from (3.1) that on passing through a weak jump in compression, i.e., for $\Delta\theta<0$, $T'/T_\infty>0$ and $\rho'/\rho_\infty>0$. Thus according to (3.2) behind a weak shock wave there is an increment in the population of the rotational quantum level with quantum number I_n such that $\Theta I_n(1+I_n)/T_\infty>1$, both on account of the rise in density and also on account of the redistribution with respect to I_n which occurs on increasing the temperature.

For the lower rotational levels [when $\Theta I_n(1+I_n)/T_\infty<1$] the displacement of the distribution with respect to I_n on increasing the temperature behind the weak shock wave leads to a reduction in population, but taken together with the rise in density the population increases. In both cases the rise in G/G_∞ due to the increase in population is greater than the fall in G/G_∞ due to the Doppler broadening of the spectral line [second term in (3.2)]. As a result of this there is a rise in G/G_∞ behind the weak shock wave for all values of the rotational quantum number. It follows from (3.3) that the larger the rotational quantum number I_n or the smaller the progressive temperature of the gas before the jump, the more significant is the increment in G/G_∞ for the same value of M_∞. For a wave of rarefaction $T'/T_\infty<0$ and $\rho'/\rho_\infty<0$ so that $G/G_\infty<1$.

In the case of purely collision broadening of the spectral line $H(a,0)/H(a_\infty,0)\approx\rho_\infty/\rho$ and Eq. (1.1) may be reduced to the form

$$\frac{G}{G_\infty}=1-\left(\frac{1}{2}\right)\frac{T'}{T_\infty}+\left[\frac{\Theta I_n(1+I_n)}{T_\infty}-1\right]\frac{T'}{T_\infty}. \tag{3.4}$$

We see from Eq. (3.4) that a rise in the translational temperature behind the weak jump in compression for low quantum levels I_n reduces the population of these levels and leads to the broadening of the spectral line by virtue of collisions. Therefore behind the weak jump there is a decrease in the value G/G_∞. In the case of $\Theta I_n(1+I_n)/T_\infty>1$, i.e., the upper rotational levels, a rise in temperature behind the jump leads to an increase in their population, and for a certain rotational quantum number $I_n=I_*$ the rise in G/G_∞ accounted for by the increase in population starts predominating over the fall in G/G_∞ due to collision broadening. By using Eq. (3.1) we may reduce (3.4) to the form

$$\frac{G}{G_\infty}=1+\frac{2\varkappa M_\infty^2\Delta\theta}{\sqrt{M_\infty^2-1}\,(7+\alpha_{H_2O})}\left[\frac{3}{2}-\frac{\Theta I_n(1+I_n)}{T_\infty}\right]. \tag{3.5}$$

The expression in square brackets is greater than zero for values of $I_n<I_*=\sqrt{1/4+3/2(T_\infty/\Theta)}-1/2$ and smaller than zero for $I_n>I_*$.

Thus for $I_n>I_*$ a weak shock wave leads to a rise in the amplification factor and a wave of rarefaction to a fall. For $I_n<I_*$ the shock wave reduces the amplification factor and the wave of rarefaction increases it. For $T_\infty=300$°K and $I_*=28$ these conclusions regarding the weak shock wave agree with those drawn in [2] after numerical shock-wave calculations.

In the range of values $0.2<a<1.4$ in which the broadening of the spectral line is determined by the collision and Doppler mechanisms $H(a,0)/H(a_\infty,0)\approx(1+1.5a_\infty)/(1+1.5a)$ and the change in the amplification factor in the perturbation wave has the form

$$\frac{G}{G_\infty}=1-\left[\frac{1}{2}\frac{T'}{T_\infty}+\frac{1.5a_\infty}{1+1.5a_\infty}\frac{\rho'}{\rho_\infty}\right]+\left\{\frac{\rho'}{\rho_\infty}+\frac{T'}{T_\infty}\left[\frac{\Theta I_n(I_n+1)}{T_\infty}-1\right]\right\}. \tag{3.6}$$

The expression in the curly brackets corresponds to the change in the population of the rotational quantum level and the expression in the square brackets, to the change in spectral linewidth due both to the collisions and to the Doppler effect. A weak jump in compression produces a fall in G/G_∞ on account of the broadening of the spectral line (which is determined by the density-increment factor) and the temperature at the leading edge of the jump, while a rise in G/G_∞ behind the weak front occurs as a result of the increase in population. Using (3.1) we may transform (3.6) to the form

$$\frac{G}{G_\infty} = 1 + \frac{M_\infty^2 \Delta\theta}{\sqrt{M_\infty^2 - 1}} \left\{ \frac{2}{5 + \alpha_{H_2O}} \left[\frac{3}{2} - \frac{\Theta I_n (1 + I_n)}{T_\infty} \right] - \frac{1}{1 + 1.5 a_\infty} \right\}. \tag{3.7}$$

It follows from this that for

$$I_n > I_{**} = -\frac{1}{2} + \sqrt{\frac{1}{4} + \frac{T_\infty}{\Theta}\left[\frac{3}{2} - \frac{5 + \alpha_{H_2O}}{2(1 + 1.5 a_\infty)}\right]}$$

there is an increase in the amplification coefficient in a shock wave and a diminution in a wave of rarefaction. If $I_n < I_{**}$, the amplification factor is reduced in the shock wave and increased in the wave of rarefaction.

We should note the way in which the change in amplification factor depends on the number M_∞. For $M_\infty = \sqrt{2}$, we see from Eqs. (3.5)-(3.7) that an extremum of the function is reached; this is associated with the behavior of the gasdynamic parameters close to $M = 1$ [7].

By considering the examples of flow around an obtuse angle or around a sharp wedge we see that weak perturbations of the gasdynamic parameters lead to a change in the amplification factor of a weak signal for the P branch of the vibrational–rotational (00^01)- (10^00) transition of the CO_2 molecules in the sense of a rise or fall, according to the particular range of values of the number a and the rotational quantum number I_n.

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