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RESONANT OSCILLATIONS OF A GAS IN AN OPEN-ENDED TUBE IN A WEAK TURBULENCE REGIME

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An analytical theory of resonant oscillations of a gas in an open-ended tube is developed. The gas flow in the tube is assumed to be turbulent. A model of gas flow near the open end of the tube is constructed. This model allows a boundary condition that is free of empirical parameters to be obtained. Theoretical results are in reasonable agreement with experimental data obtained by other authors.

The theory of resonant oscillations of a gas in tubes-resonators is one of the most interesting problems of hydrodynamic acoustics. High-intensity oscillations of a gas are usually excited by a piston which harmonically oscillates at one end of the tube [1–4]. Of particular interest from the practical viewpoint are open-ended tubes-resonators. Oscillations in such systems are accompanied by a number of interesting effects: an oscillating jet is formed at the open end of a tube [5], a nonuniform temperature field is established in a tube [6], etc.

The qualitative theory of the phenomenon is not yet complete. This is connected with the complexity of the boundary condition at the open end of a tube [7, 8] and poor knowledge of the specifics of an oscillating turbulent flow in a tube [9]. An analytical model of the processes at the open end of the tube has been constructed recently, and, using this model, the boundary condition has been determined [10]. The models of tube turbulence, which were proposed in [11, 12] for the first time, offer a description of experimental results, but have significant drawbacks: (a) they ignore the heat transfer between the tube wall and a gas; (b) they are based on the assumption of a quasi-stationary regime of turbulence; (c) they do not consider dispersion in a turbulent medium.

In this paper, an attempt is made to construct a model of resonant oscillations of a gas at one end of a tube in a turbulent flow regime that is free of the above drawbacks.

Oscillations in a cylindrical tube of length L and radius R, which are excited by a harmonically oscillating piston with displacement amplitude $l_0 \ll L$, are characterized by the following dimensionless parameters [7–9, 13]:

$$\varepsilon = \frac{V}{\omega L}, \quad H = R\sqrt{\frac{\omega}{\nu}}, \quad \mathrm{sinh} = \frac{\omega R}{V}, \quad \mathrm{M}_p = \frac{\omega l_0}{c_0}, \quad \mathrm{Re}_\omega = \frac{V^2}{\omega \nu}.$$

Here V is the amplitude of velocity fluctuations in a velocity loop (for the first resonance, at the open end of the tube), ω is the cyclic frequency of oscillations, ν is the kinematic viscosity, and c_0 is the speed of sound in an undisturbed gas. Since $l_0 \ll L$, for oscillations around the basic resonance frequency $\omega_0 = \pi c_0/(2L)$ [13] we obtain $\mathrm{M}_p \ll 1$. In experiment, we usually have $H \gg 1$ (the effect of the acoustic boundary layer on the flow core is small) and $\mathrm{sinh} \leqslant 1$. For a long tube ($L/R \gg 1$), the condition $\mathrm{sinh} \leqslant 1$ leads to $\varepsilon \ll 1$, i.e., the problem can be solved by the methods of disturbance theory [13]. The criterion Re_ω indicates a turbulence regime: if $10^5 \leqslant \mathrm{Re}_\omega \leqslant 6 \cdot 10^5$, the regime of weak turbulence occurs [9]. This regime is of interest because the majority of experiments were conducted under these conditions [1–4].

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Nonisentropic oscillations of a gas in a tube can be described by the equations [14]

$$\frac{\partial(\rho u)_s}{\partial t} + \frac{\partial}{\partial x}(\rho u^2)_s + \frac{\partial p}{\partial x} + \frac{2\tau}{R} = 0, \qquad \frac{\partial p}{\partial t} + u_s\frac{\partial p}{\partial t} + æp\frac{\partial u_s}{\partial x} - \frac{2(æ-1)}{R}q = 0, \tag{1}$$

where ρ, p, and u are the density, pressure, and velocity, respectively; τ and q are the shear stress and the heat flux at the tube wall, $æ = c_p/c_V$, t is the time, x is the longitudinal coordinate (the closed end of the tube corresponds to $x = 0$, and the open end to $x = L$), and the subscript s means that the quantity is averaged over the tube cross section. In the first (acoustic) approximation, we obtain from (1) that

$$\rho_0\frac{\partial u_{1s}}{\partial t} + \frac{\partial p_1}{\partial x} = -\frac{2\tau_1}{R}, \qquad \frac{\partial p_1}{\partial t} + \rho_0 c_0^2\frac{\partial u_{1s}}{\partial x} = \frac{2(æ-1)}{R}q_1. \tag{2}$$

The subscript 1 denotes here the first approximation, and the subscript 0 refers to the quantities of an undisturbed flow. To solve system (2), one must estimate τ_1 and q_1.

First of all, we take into account that with allowance for the boundary condition $u(R) = 0$, the Reynolds equations written near the tube wall lead to the relation

$$\frac{\partial p}{\partial x} = -\frac{1}{r}\frac{\partial}{\partial r}(r\tau_1)\Big|_{r=R}.$$

It follows from here that, since p is independent of r, the dependence of τ_1 on x and time is similar to the dependence of the pressure gradient on x. On the other hand, the Reynolds equations written at the tube axis with allowance for the symmetry condition

$$\frac{1}{r}\frac{\partial}{\partial r}(r\tau_1)\Big|_{r=0} = 0$$

yield

$$\rho_0\frac{\partial u_1}{\partial t}\Big|_{r=0} + \frac{\partial p_1}{\partial x} = 0,$$

i.e., the dependence of the pressure gradient on x is similar to that of the velocity. Summarizing the aforesaid, we can show that in (2)

$$\tau_1 \propto u_1(x, t + \phi). \tag{3}$$

The validity of (3) is verified by experiment [9]. For a uniform velocity distribution along the tube, the relationship between the shear-stress amplitude on the wall τ_{1m} and the amplitude of velocity fluctuations at the tube axis is as follows [9]:

$$\tau_{1m} = \frac{1}{2}\rho_0 f_w (u_{1m})^2. \tag{4}$$

In the case of weak turbulence, the friction coefficient on the wall is $f_w \approx 0.005$ [9].

For (3) to be consistent with (4), we assume that

$$\tau_1(x,t) = \rho_0\beta_0 u_1(x, t+\phi), \qquad \beta_0 = \frac{Vf_w}{2}\int_0^L \frac{u_{1m}(x)}{V}\,dx, \tag{5}$$

where the coefficient β_0 takes into account that the amplitude of velocity fluctuations u_{1m} depends on x for resonant oscillations of a gas in a tube.

Before using (5) in (2), we pass from the flow velocity at the tube axis u_1 to the cross section-averaged velocity u_{1s} and take into account the phase shift between the oscillations τ_1 at the wall and u_1 at the axis. In a weak turbulence regime, the amplitude profile of velocity fluctuations is uniform everywhere, except for a layer of thickness δ_1, in which a universal distribution of the amplitude of velocity fluctuations is observed, i.e.,

$$\frac{u_{1m}}{u^*} = 2.5\ln\left(\frac{R-r}{\nu}u^*\right) + 5. \tag{6}$$

Here $u^* = (\tau_{1m}/\rho_0)^{1/2}$ is the dynamic velocity and r is the radial coordinate [9]. As follows from experimental data of [9], the layer thickness δ_1 can be found from the formula

$$\frac{\delta_1}{R} = \frac{0.0154}{\sinh}. \tag{7}$$

Using (6) and (7), we obtain

$$u_{1s} = u_{1m}B, \qquad B = 1 - (C \ln \mathrm{Re}_\omega + D)/\sinh, \tag{8}$$

where $C \approx -0.00385$ and $D \approx 0.0546$.

The empirical formula for the phase shift ϕ between τ_1 at the wall and u_1 at the tube axis is obtained from the data of [9]:

$$\phi = 0.838 - 0.891(\mathrm{Re}_\omega \cdot 10^{-6}). \tag{9}$$

With allowance for (8) and (9), we have

$$\tau_1 = \rho_0 \beta \exp(i\phi) u_{1s}, \qquad \beta = \beta_0/B. \tag{10}$$

To estimate q_1, we assume

$$q_1 = -\beta_T p_1. \tag{11}$$

The thicknesses of the dynamic and thermal boundary layers can be calculated from the formulas

$$\delta_1 = \sqrt{2\mu_e/\rho_0\omega}, \qquad \delta_{T1} = \sqrt{2\lambda_e/\rho_0 c_p \omega} \tag{12}$$

(the subscript e refers to the effective value). Since $\delta_1 \ll R$ and $\delta_{T1} \ll R$, in the definitions

$$q_1 = \lambda_e \frac{\partial T_1}{\partial r}\bigg|_w, \qquad \tau_1 = -\mu_e \frac{\partial u_1}{\partial r}\bigg|_w$$

the derivatives with respect to r can be replaced by the increment ratios. From (10) and (11), and the conditions at the tube wall $u_1(R) = 0$ and $T_1(R) = 0$, it follows that

$$\frac{\lambda_e T_{1m}}{\delta_{T1}} \approx \beta_T p_1, \qquad \frac{\mu_e}{\delta_1} \approx \rho_0 \beta \exp(i\phi). \tag{13}$$

One can easily show that $p_1 \approx \rho_0 c_p T_{1m}$ outside the boundary layer. From (12) and (13), we obtain

$$\beta_T = \beta \exp(i\phi)/\sqrt{\mathrm{Pr}_t}. \tag{14}$$

The turbulent Prandtl number is $\mathrm{Pr}_t \approx 0.9$ in the boundary layer of steady turbulent flows [15]. We assume this value to be acceptable for our case, too.

We pass to the dimensionless variables in (2), assuming that $\bar{p}_1 = p_1/\rho_0 c_0^2$ and $\bar{u}_{1s} = u_{1s}/c_0$. Taking into account (10), (11), and (14), we obtain

$$\frac{1}{c_0}\frac{\partial \bar{u}_{1s}}{\partial t} + \frac{\partial \bar{p}_1}{\partial x} = -a\bar{u}_{1s}, \quad \frac{1}{c_0}\frac{\partial \bar{p}_1}{\partial t} + \frac{\partial \bar{u}_{1s}}{\partial x} = -\frac{æ - 1}{\sqrt{\mathrm{Pr}_t}} a\bar{p}_1, \quad a = \frac{2\beta \exp(i\phi)}{Rc_0}. \tag{15}$$

The solutions of system (15) have the form

$$\begin{gathered} \bar{p}_1 = r_1 \cos(kx + \alpha_1 + i\gamma_1) \exp[i(\omega t + \psi_1)], \\ \bar{u}_{1s} = -ir_1\mu_1 \sin(kx + \alpha_1 + i\gamma_1) \exp[i(\omega t + \varphi + \psi_1)]. \end{gathered} \tag{16}$$

Here r_1, α_1, γ_1, and ψ_1 are real constants of integration, $\mu_1 = |k/(k_0 - ia)|$, $k_0 = \omega/c_0$, $\varphi = \arg[k/(k_0 - ia)]$, and

$$k^2 = k_0^2\left[1 - i\frac{a}{k_0}\left(1 + \frac{æ - 1}{\sqrt{\mathrm{Pr}_t}}\right) - \frac{æ - 1}{\sqrt{\mathrm{Pr}_t}}\frac{a^2}{k_0^2}\right]. \tag{17}$$

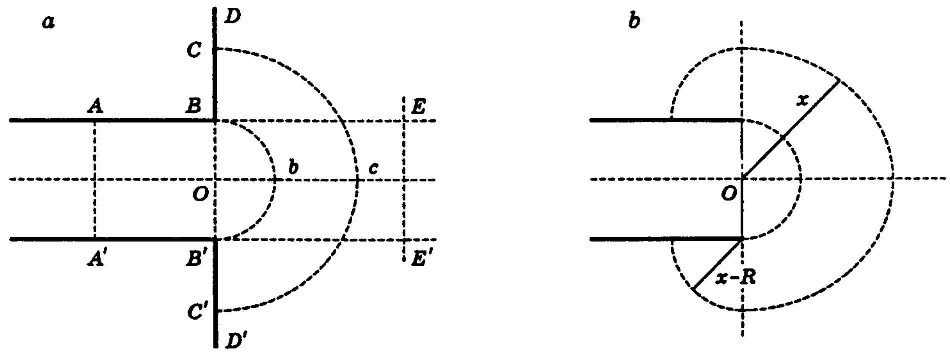

Fig. 1

Since $a/k_0 \ll 1$, the last term in (17) can be ignored, i.e.,

$$k \approx k_0 - ib, \qquad b = \frac{a}{2}\left(1 + \frac{æ - 1}{\sqrt{\mathrm{Pr}_t}}\right) = b_1 + ib_2, \tag{18}$$

where $b_1 = (\beta/Rc_0)(1+(æ-1)/\sqrt{\mathrm{Pr}_t})\cos\phi$ and $b_2 = (\beta/Rc_0)(1+(æ-1)/\sqrt{\mathrm{Pr}_t})\sin\phi$. Thus, $k \approx k_0+b_2-ib_1$; b_1 and b_2 refer to absorption and dispersion, respectively.

We now consider the boundary conditions. We prescribe the piston velocity at the tube end closed by the piston ($x = 0$):

$$\bar{u}_{1s}(0,t) = -\mathrm{M}_p \exp(i\omega t). \tag{19}$$

We consider the flow at the open end ($x = L$) of the tube equipped with an infinite flange (Fig. 1a). Let, at a certain distance from the exit cross section inside the tube (cross section AA'), the velocity variation obey the law

$$u = V\cos\omega t. \tag{20}$$

We use a model proposed in [16, 17] which implies the jet outflow ($u \geqslant 0$) and the spherical inflow ($u \leqslant 0$) into an orifice located at the point O. The gas leaving the tube is enclosed in a volume with generatrices BE and BE'. Since the mixing layer of the jet has no time to develop near the open end of the tube (at distances $x < 6R$ [5]), the jet cross-sectional area remains practically constant and equal to the tube cross-sectional area S_0. In this case, the gas velocity is also independent of x. The gas flows into the tube through the hemispheres BbB', CcC', etc. The role of viscous losses in the source of suction is insignificant [18]. Thus, we can assume that the gas inflow is potential and the hemispheres are isotachs. The amount of gas crossing the hemispheres remains constant, and the following equation is valid for a hemisphere of arbitrary radius x:

$$u(x,t) = \Phi(x)u_{BbB'}(t) \qquad [\Phi(x) = R^2/x^2], \tag{21}$$

where $u_{BbB'}(t)$ is the velocity at the points of the hemisphere BbB' passing through the tube edges ($x = R$). For a tube without a flange (Fig. 1b), we have $\Phi(x) = R^2/[x^2 + (x-R)^2 + (\pi/2)R(x-R)]$.

We consider the outflow through the cross section BB' and the inflow through the hemisphere BbB'. By virtue of the mass conservation law, the amount of gas going out through BB' should be compensated by a return gas flow through BbB', i.e.,

$$S_0\int_0^{t_1} u_{1s}(t)\,dt + S\int_{t_1}^{T} u_{BbB'}(t)\,dt = 0 \tag{22}$$

($S = 2\pi R^2$). Since $S > S_0$, to satisfy (22) it is necessary that the outflow duration t_1 be larger than the inflow duration. This is possible if the velocity has a constant component. Assuming the latter to be proportional to the amplitude of fluctuations of the velocity V, we obtain [17] the expression $u = V(m_0 + \cos\omega t)$, $x = R$, where the parameter m_0 should be determined from (22).

The outflow duration t_1 is found from the condition $u = 0$. Then we have

$$\begin{aligned} u_{1s}(t) &= BV(m_0 + \cos\omega t), \qquad -(\pi/2+\theta) \leqslant \omega t \leqslant (\pi/2+\theta), \\ u_{BbB'}(t) &= V(m_0 + \cos\omega t), \qquad (\pi/2+\theta) \leqslant \omega t \leqslant (3\pi/2-\theta), \end{aligned} \tag{23}$$

where $\theta = \arcsin m_0$. Substituting (23) into (22), we obtain

$$(B+2)\pi m_0 + 2(B-2)[m_0 \arcsin m_0 + \sin(\arccos m_0)] = 0. \tag{24}$$

The calculation shows that $B \approx 0.93$ under the experimental conditions of [3]. It follows from (24) that $m_0 \approx 0.239$.

We consider the oscillations of the particles intersecting, for instance, the cross section EE' (Fig. 1a), assuming the motion to be potential. The outflow velocity is determined by (23), and the inflow velocity is found from (21) according to which the velocity decreases rapidly as x increases. Expanding $u(t)$ into a Fourier series, we have

$$\begin{gathered} \bar{u} = \mathrm{M}_E\Big\{\Big(\frac{m_0}{2+a_0}\Big) + \Big(\frac{m_0}{2-a_0}\Big)\Phi(x) + \Big[\Big(\frac{1}{2}+a_1\Big) + \Big(\frac{1}{2}-a_1\Big)\Phi(x)\Big]\cos\omega t + a_2[1-\Phi(x)]\cos 2\omega t + \ldots\Big\}, \\ a_0 = \frac{1}{\pi}(m_0\theta + \cos\theta), \qquad a_1 = \frac{1}{\pi}\Big(\theta + 2m_0\cos\theta - \frac{1}{2}\sin 2\theta\Big), \\ a_2 = \frac{1}{\pi}\Big(\cos\theta - m_0\sin 2\theta - \frac{1}{3}\cos 3\theta\Big), \qquad \mathrm{M}_E = \frac{V}{c_0}, \qquad \bar{u} = \frac{u}{c_0}. \end{gathered} \tag{25}$$

An analysis shows that significant changes in $\bar{u}$ with distance x, which are faster in the case of a tube without a flange, end at $x \approx 3R$. Beginning from $x \approx 5R$, the composition of oscillations becomes independent of x and the geometry of the tube's open end. For $x > 6R$, the viscosity becomes significant. The evolution of an oscillating jet was considered in detail in [5]. Let $\bar{u}_\infty$ be the velocity at a large distance from the open end of the tube, where we can still ignore viscous effects (let this be valid for the cross section EE'). Assuming $x \to \infty$ in (25), we have

$$\bar{u}_\infty = \mathrm{M}_E[(m_0/2 + a_0) + (1/2 + a_1)\cos\omega t + a_2\cos 2\omega t + \ldots]. \tag{26}$$

For a potential flow whose velocity is described by (25), we can use the Lagrange–Cauchy integral

$$\frac{p}{\rho_0} + \frac{u^2}{2} + \frac{\partial\varphi}{\partial t} = F(t), \tag{27}$$

where the order of the third term is estimated as Sh, i.e., it can be ignored. Applying (27) to two cross sections (for example, AA' and EE' in Fig. 1a), assuming the atmospheric pressure in the cross section EE', and obtaining the velocity from (20) and (27), after simple transformations we have

$$\bar{p}_1(L,t) = m\bar{u}^0_{1s}\bar{u}_{1s}(L,t), \quad m = m_1/B^2, \quad m_1 = (1/2 + a_1)(m_0/2 + a_0 + a_2/2), \tag{28}$$

where $\bar{u}^0_{1s}$ is the amplitude of velocity fluctuations averaged over the tube cross section at the open end of the tube. Under the experimental conditions of [3], we have $m_1 \approx 0.361$.

Condition (28) is nonlinear, as in [3, 6–8, 11], but, in contrast to the variants used in these papers, it is *derived* from the flow model near the open end of the tube without any semiempirical considerations.

Substituting the solutions of (16) into (19) and (28), we obtain a system for determination of the desired constants:

$$\begin{gathered} r_1\mu_1 \sin\alpha_1 \cosh\gamma_1 = \mathrm{M}_p\cos(\varphi + \psi_1), \quad r_1\mu_1\cos\alpha_1\sinh\gamma_1 = -\mathrm{M}_p\sin(\varphi+\psi_1), \\ \cos z\cosh w = mr_1\mu_1^2\sqrt{\sin^2 z + \sinh^2 w}(\cos z\sinh w\cos\varphi + \sin z\cosh w\sin\varphi), \\ \sin z\sinh w = mr_1\eta_1^2\sqrt{\sin^2 z + \sinh^2 w}(\sin z\cosh w\cos\varphi - \cos z\sinh w\sin\varphi). \end{gathered} \tag{29}$$

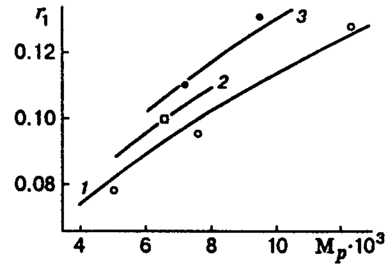

Fig. 2

TABLE 1

$M_p \cdot 10^3$	L/R	$\Delta\omega/\omega_0$	
		Theory	Experiment
5	171	0.020	0.025
7.6	171	0.022	0.033
12,35	171	0.023	0.030
6.6	129	0.021	0.027
7.2	89	0.022	0.023
9.5	89	0.022	0.027

Here $z = (k_0 + b_2)L + \alpha_1$ and $w = \gamma_1 - b_1 L$. System (29) is easily solved under the assumptions $r_1 \ll 1$, $\sinh w \sim r_1$, $\cosh w \sim 1$, and $\mu_1 \sim 1$, whence $\varphi_1 \ll 1$. We write the solution in the form

$$\alpha_1 = \frac{\pi}{2} - (k_0 + b_2)L, \quad \gamma_1 = mr_1 + b_1 L, \quad \sin(\varphi + \psi_1) = -\frac{r_1(mr_1 + b_1 L)\sin(k_0 + b_2)L}{M_p},$$
$$r_1\sqrt{\cos^2(k_0 + b_2)L + (mr_1 + b_1 L)^2 \sin^2(k_0 + b_2)L} = M_p. \tag{30}$$

It follows from the above equations that the resonance in the system is reached when $(k_0 + b_2)L = \pi/2$, i.e., the resonance frequency is shifted, and, for this shift, from formulas (5), (10), and (18) we obtain

$$\frac{\pi/2 - k_0 L}{\pi/2} = \frac{2}{\pi} b_2 L = \frac{2}{\pi}\frac{f_w r_1}{3B}\frac{L}{R}\left(1 + \frac{æ - 1}{\sqrt{\mathrm{Pr}_t}}\right)\sin\phi. \tag{31}$$

Under the resonance conditions, it follows from (30) that

$$r_1 = M_p^{1/2}\left[m + \frac{f_w}{3B}\frac{L}{R}\left(1 + \frac{æ - 1}{\sqrt{\mathrm{Pr}_t}}\right)\cos\phi\right]^{-1/2}, \tag{32}$$

which again takes into account (5), (10), and (18).

Figure 2 shows experimental data [3] (points) and calculation results obtained by (32) (curves). Agreement of the data is quite satisfactory: the points deviate from the dependence (32) by no more than 5%. The data scatter is caused, as follows from (32), by different L/R ratios ($L/R = 171$, 129, and 89 for curves 1–3, respectively).

We now discuss the shift of the resonance frequency $\Delta\omega/\omega_0$ described by formula (31). Before comparing the values calculated from (31) with experimental data, we have to take into account the so-called tip correction to the tube length, which is induced by flow inertia near the open end [3], i.e., to substitute L for $L + \Delta R$ in (31), where $\Delta \sim 1$ [3]. For the resonance-frequency shift, we can write

$$\frac{\Delta\omega}{\omega_0} = \frac{2}{\pi}\frac{f_w r_1}{3B}\frac{L}{R}\left(1 + \frac{æ - 1}{\sqrt{\mathrm{Pr}_t}}\right)\sin\phi + \Delta\frac{R}{L}.$$

The maximum discrepancy between the calculated and experimental results is 33% (see Table 1).

Thus, we can state that the model proposed in the present work is in good agreement with the available experimental data.

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