SUPERPOSITION PRINCIPLE IN ACOUSTICS

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Using the effect of the nonstationary frequency shift of acoustic waves in a medium with spatially and temporally varying properties, a theoretical estimate is made of the deviation from the superposition principle for an acoustic wave propagating in the field of another continuous plane wave.

According to the superposition principle of waves [1], in a linear medium each free wave propagates independently of all remaining waves, and the sound field at each point is simply the sum of fields of the component free waves. For scalar characteristics of the wave (e.g., for pressure, temperature) the summation is algebraic, and for vector properties (velocity, particle acceleration) it is vector summation.

The superposition is approximate: it is valid to the extent that linearization of the equations of hydrodynamics is satisfied for sound waves.

The error due to linearization is small in the same sense that the terms neglected in the equations are small in comparison with the terms retained. This does not imply, however, that the error remains small in the solution of the equation at all times of the wave motion. On the contrary, the error due to linearization accumulates with the wave propagation: the further the wave propagates, the more its profile is deformed. There exists here a definite analogy with wave damping. Forces of internal friction in a sound wave are negligbly small in comparison with elastic forces. If they are not taken into account, however, damping does not occur.

It must be noted that although in theoretical acoustics the superposition principle is correctly treated, most experimenters and practioners, associated with metrological uses of acoustics, particularly ultrasound, perceive this principle as absolute truth. A measurement of the deviation from the superposition principle can provide useful information on its properties. In other cases the deviation from the superposition principle can serve as an additional source of error in acoustic measurements. If some probing acoustic signal propagates in a medium perturbed by extraneous acoustic sources, during its propagation it will accumulate information on the parameters of the extraneous sources, which can also be used in practice.

Using the theory of nonstationary Doppler effect developed in [2-4], the relations obtained make it possible to estimate the deviations from the superposition principle. The medium, disturbed by the external acoustic field, can be treated as a medium whose properties vary in space and time, and, as shown in the papers mentioned, the propagation of acoustic waves in it is accompanied by a change of their frequency. There are two reasons leading to a frequency shift in this case. The first reason is a change of pressure in the medium due to the external acoustic perturbations. In the case of a liquid, for example, the velocity of sound is determined by the equation

$$v = \sqrt{1/\rho\beta}. \tag{1}$$

For small perturbations we put

$$\rho = \rho_0 (1 + k_1 \Delta p), \ \beta = \beta_0 (1 + k_2 \Delta p), \tag{2}$$

where ρ_0 and β_0 are the density and compressibility of the unperturbed medium, and Δp is the excess (sound) pressure due to the action of the external source of acoustic field. Substituting relation (2) into (1), we obtain

$$v = v_0 / \sqrt{1 + (k_1 + k_2) \Delta p + k_1 k_2 \Delta p^2}, \tag{3}$$

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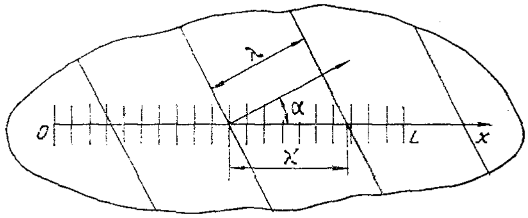

Fig. 1. Relative position of acoustic waves.

whence it is seen that in real media the velocity of sound always depends on pressure, even when the condition $k_1 = -k_2$ is satisfied, where $v_0 = 1\sqrt{\rho_0\beta_0}$ is the velocity of sound in the medium unperturbed by the external acoustic field. The second reason leading to violation of the superposition principle consists of the fact that the action of the external acoustic field generates vibrational motion of the medium particles, and if probing elastic waves pass through this medium, the frequency of the latter changes due to the convective Doppler effect [4]. Thus, if elastic waves pass through a medium perturbed by an external source of an acoustic field, at the exit from the medium the waves mentioned are frequency or phase modulated by the external sound field.

We estimate the frequency—phase shift generated by an external acoustic field in elastic waves. In this case we consider separately the effect of sound pressure and the effect of accelerated motion of the medium, with the purpose of simplifying the problem.

Let continuous planar longitudinal waves propagate in the direction of the x axis. For definiteness we assume that their source is located at the point x = 0, and that the receiver is at x = L (see Fig. 1); the frequency of the emitted waves is constant and equals f_0. At the same time there exists and external acoustic field in the medium: the continuous plane waves with frequency F propagate in some direction, forming an angle α with the x axis. For clarity we assume that $F \ll f_0$. The propagation velocity of the probing wave (of frequency f_0) at the moment of its emission at point x = 0 will depend on the amplitude and phase of the external acoustic field at this point, which determine the sound pressure in the medium, the temperature variation, etc. In the case of water, e.g., for $\Delta p > 0$ the propagation velocity of the probing wave is larger, and for $\Delta p < 0$ smaller than its propagation velocity in the unperturbed medium. This leads to the consequence that the propagation time of the probing wave from the emitter to the receiver (the delay time τ_d) will depend on the initial conditions, the amplitude and phase of the external field at point x = 0 at the moment of emission of the probing wave. Since the initial conditions change periodically, the delay time is also a periodic function of time. This leads to a periodic frequency shift of the applied probing waves: at the receiver there will be a frequency-modulated (or phase-modulated) signal, while the modulation frequency equals F.

As seen from Fig. 1, the following relation is valid

$$\lambda' = \lambda/\cos\alpha, \tag{4}$$

and for the propagation velocity of the probing waves in their direction of propagation one can write:

$$v(x, \tau) = v_0 + \Delta v_0 \sin\left(\Omega\tau - \frac{2\pi}{\lambda'}x + \varphi\right) = v_0 + \Delta v_0 \sin\Omega\left(\tau - \frac{x\cos\alpha}{v_0} + t\right), \tag{5}$$

where $\Omega = 2\pi F$; $t = \varphi/\Omega$.

The essence of the method suggested below for determining the frequency shift and finding the law of wave motion in the medium with account of the space—time dependence of its propagation velocity and subsequently finding the relation between the frequency of the emitted wave with the frequency that this wave has upon approaching the receiver. It is clear from physical considerations that the largest frequency shift will be observed at $\alpha = 0$, when probing an external acoustic waves propagate in the same direction. In this case any phase surface of the probing wave has plenty of time of being found either in the region of enhanced or reduced pressure, and, correspondingly, will gradually either lead the external wave, or lag behind it. For $\alpha \neq 0$ the probing wave will subsequently be in the region of enhanced and depressed pressure, and the frequency shift accumulated during its propagation in the region of enhanced pressure will, to some extent, compensate the inverse sign of frequency shift accumulated in the region of reduced pressure. Therefore, we confine ourselves to the case of $\alpha = 0$.

To determine the law of motion of a wave emitted in the direction of the x axis at some moment of time, it is necessary to solve the differential equation

$$\frac{dx}{d\tau} = v_0 + \Delta v_0 \sin \Omega (\tau - x/v_0 + t) \tag{6}$$

with corresponding initial conditions. Without loss of generality, the initial condition can be taken in the form

$$x|_{\tau=0} = 0. \tag{7}$$

The solution of Eq. (6) satisfying condition (7) is

$$\ln \left| \frac{\mathrm{tg} \frac{\Omega}{2} \left(\tau - \frac{x}{v_0} + t \right)}{\mathrm{tg} \frac{\Omega}{2} t} \right| = - \frac{\Omega \Delta v_0 \tau}{v_0}. \tag{8}$$

The latter expression is the law of motion of the wave emitted at point $x = 0$ at moment of time $\tau = 0$. If the running coordinate x is replaced by the coordinate $x = L$, at which the receiver is placed, the time found from this expression is the delay time of the wave path from the source to the receiver, i.e., the delay time τ_d. Thus,

$$\ln \left| \frac{\mathrm{tg} \frac{\Omega}{2} \left(\tau_d - \frac{L}{v_0} + t \right)}{\mathrm{tg} \frac{\Omega}{2} t} \right| + \frac{\Omega \Delta v_0}{v_0} \tau_d = 0. \tag{9}$$

For nonvarying L, Δv_0, v_0, Ω the quantity τ_d is a function of t, where t can be considered as a running time relative to the radiator. Indeed, the phase φ of the traveling external waves at the position of the emitter varies continuously and linearly with time, $\varphi = \Omega t$, where Ω = const. Therefore, in what follows we replace t by τ for the running time. Further, it is seen from Eq. (6) that for $t = \pi n/\Omega$; $n = 0, 1, 2, \ldots$, which corresponds to the initial phase $\varphi = \pi n$, its solution is the dependence $x = v_0 \tau$, i.e.,

$$\tau_d^0 = L/v_0. \tag{10}$$

This implies that the delay time for other values of φ cannot differ from τ_d^0 by more than π/Ω. Therefore, for convenience of further analysis the quantity τ_d is replaced by the quantity equal to it

$$\tau_d = \frac{L}{v_0} + \Delta \tau_d. \tag{11}$$

Substituting (11) into Eq. (9) and replacing t by τ, we obtain a transcendental equation for determining the quantity $\Delta\tau_d$ as a function of time:

$$\ln \left| \frac{\mathrm{tg} \frac{\Omega}{2} (\Delta \tau_d + \tau)}{\mathrm{tg} \frac{\Omega}{2} \tau} \right| + \left(\frac{L}{v_0} + \Delta \tau_d \right) \frac{\Omega \Delta v_0}{v_0} = 0. \tag{12}$$

To determine the frequency shift we write the obvious relation for the phase difference of the received φ_r and emitted φ_e oscillations:

$$\varphi_r - \varphi_e = \Delta \varphi = - 2\pi f_0 \tau_d. \tag{13}$$

Since the quantity τ_d varies in time, so does the phase difference mentioned, with resulting equivalent frequency shift

$$\Delta f = \frac{1}{2\pi} \frac{d(\Delta\varphi)}{d\tau} = - f_0 \frac{d\tau_d}{d\tau} = - f_0 \frac{d(\Delta\tau_d)}{d\tau}. \tag{14}$$

The derivative $d(\Delta\tau_d)/d\tau$ can be found as a derivative of an implicit function, which is the left-hand side of Eq. (12). Differentiating and substituting into (14), we obtain the following equation for the frequency shift:

$$\Delta f(\tau) = -f_0 \frac{\sin\Omega(\Delta\tau_d - \tau) - \sin\Omega\tau}{\left[1 + \frac{\Delta v_0}{v_0}\sin\Omega(\Delta\tau_d + \tau)\right]\sin\Omega\tau}. \tag{15}$$

Since the condition $\Delta v_0/v_0 \ll 1$ is practically always satisfied, the following expression can be used quite accurately

$$\Delta f(\tau) \approx f_0\left[1 - \frac{\sin\Omega(\Delta\tau_d + \tau)}{\sin\Omega\tau}\right]. \tag{16}$$

Taking into account that for $\varphi \to \pi n \Delta\tau_d \to 0$, it is easily verified that the function $\Delta f(\tau)$, described by expression (16), achieves a maximum at $\varphi = 2\pi n$, which corresponds to the values $\tau = 2\pi n/\Omega$, and minima at $\varphi = (2n + 1)\pi$, which corresponds to $\tau = (2n + 1)\pi/\Omega$. Physically this is explained by the fact that at a given moment of time there exists, respectively, at the position of the emitter a maximum velocity of increasing or decreasing sound pressure from the traveling external wave, which also generates the largest frequency shift.

Consider, as an example, the case $\varphi \to 0$ $(\tau \to 0)$. It follows from the condition $\lim_{\tau\to 0}\Delta\tau_d \to 0$ that for $\tau \to 0$ we can put

$$\Delta\tau_d = a\tau, \tag{17}$$

where a is some number. Substituting the latter into Eq. (12) and passing to the limit $\tau \to 0$, we obtain

$$\ln|a + 1| = -\frac{\Omega\Delta v_0 L}{v_0^2}$$

or

$$|a + 1| = \exp\left[-\frac{\Omega\Delta v_0 L}{v_0^2}\right].$$

The right-hand side of the last equality is positive and smaller than unity; hence, it can be concluded that $a < 0$, and $|a| < 1$. In this case the absolute value sign can be neglected, and, thus,

$$a = \exp\left[-\frac{\Omega\Delta v_0 L}{v_0^2}\right] - 1. \tag{18}$$

Similarly we find now the limiting expression of (16) with account of (17) and (18):

$$\lim_{\tau\to 0}\Delta f(\tau) = -f_0 a = f_0\left(1 - \exp\left[-\frac{\Omega\Delta v_0 L}{v_0^2}\right]\right). \tag{19}$$

Under realistic conditions the condition $\Omega\Delta v_0 L/v_0^2 \ll 1$ is usually satisfied; therefore, it can be assumed that quite accurately

$$\Delta f \approx f_0 \frac{\Omega\Delta v_0 L}{v_0^2}. \tag{20}$$

For the initial phase $\varphi = (2n + 1)\pi$ we obtain similarly

$$\Delta f = -f_0\left(\exp\left[\frac{\Omega\Delta v_0 L}{v_0^2}\right] - 1\right) \tag{21}$$

or approximately

$$\Delta f \approx -f_0 \frac{\Omega\Delta v_0 L}{v_0^2}. \tag{22}$$

It is seen from relations (19), (21) or (20), (22) that the frequency shift increases with the distance from the source, and the error due to the deviation from the superposition principle indeed accumulates as the wave propagates. These relations also make it possible to carry out a numerical estimate of the maximum frequency shift generated under the conditions of the problem considered. The estimate is performed for liquid (water) and gas (air) media. To determine Δv_0 we take into account that $\Delta v_0 \ll v_0$, $\Delta p_0 \ll p_0$, and one can start from the relation

$$\Delta v_0 = \left(\frac{\partial v}{\partial p}\right)_T \Delta p_0 + \left(\frac{\partial v}{\partial T}\right)_p \Delta T_0. \tag{23}$$

According to the data of [5], for water at normal pressure $\partial v/\partial p = 1.8 \cdot 10^{-6}$ m/sec·Pa; for ΔT_0 in [6] the following relation is suggested

$$\frac{\Delta T_0}{T_0} \approx \frac{\alpha_T v_0^2}{c_p}\left(\frac{u_0}{v}\right) = \frac{\alpha_T \Delta p_0}{\rho_0 c_p}.$$

Taking this into account, instead of (23) one can write

$$\Delta v_0 = \left[\left(\frac{\partial v}{\partial p}\right)_T + \frac{\alpha_T T_0}{\rho_0 c_p}\left(\frac{\partial v}{\partial T}\right)_p\right] \Delta p_0. \tag{24}$$

We use the following values: $f_0 = 10^6$ Hz, $\Omega = 2\pi \cdot 10^5$ rad/sec, L = 10 m, v_0 = 1500 m/sec, T_0 = 300°K, Δp_0 = 1 Pa. Substituting in Eq. (24) the values for water $\alpha_T = 2 \cdot 10^{-4}$ K^{-1}, $\rho_0 = 10^3$ kg/m^3, $c_p = 4.2 \cdot 10^3$ J/kg·°K, dv/dT = 3 m/sec, we obtain $\Delta v_0 = 1.84 \cdot 10^{-6}$ m/sec. It seems in this case that the velocity increase due to hydrostatic pressure is $1.8 \cdot 10^{-6}$ m/sec, and due to adiabatic heating of water in the sound wave it is $4.2 \cdot 10^{-8}$ m/sec, i.e., a negligibly small quantity. Substituting these data into Eq. (22), we find the value of the maximum frequency shift $\Delta f \approx 5$ Hz. For L = 50 m the frequency shift is 25 Hz. The resolving capability of available receivers of frequency-modulated signals makes it possible to measure these frequency shifts.

For air, taking into account the large attenuation coefficient of ultrasound at high frequencies, we adopt other conditions: $f_0 = 10^5$ Hz, $\Omega = 2\pi \cdot 10^4$ rad/sec, L = 1 m. According to experimental data for air, dv/dp = $1.3 \cdot 10^{-6}$ m/sec·Pa [7]; one easily calculates d T = 0.6 m/sec·°K. The temperature change in the external wave can be determined from the equation [8]

$$\Delta T_0 \approx \frac{\gamma - 1}{\gamma} \frac{\Delta p_0}{p_0} T_0, \tag{25}$$

Where $\gamma = 1.4$; for T_0 = 300°K, $p_0 = 10^5$ Pa, Δp_0 = 1 Pa we obtain $\Delta T_0 = 8.6 \cdot 10^{-4}$ K. Substituting these values into Eq. (23), we obtain for air $\Delta v_0 = 5.2 \cdot 10^{-4}$ m/sec; from Eq. (20) the frequency shift equals (for v_0 = 340 m/sec) $\Delta f \approx 27$ Hz. For L = 5 m Δf = 135 Hz.

It must be pointed out that for gas media the main contribution to the velocity increment Δv_0 is due to the temperature increment in the sound wave. Under the conditions adopted for air the first term in expression (23) is approximately 400 times smaller than the second.

We estimate now the frequency shift generated by the vibrational motion of medium particles in an external acoustic field.

We use the relations for the longitudinal and transverse convective Doppler effects, obtained in [4]. Since the frequency shift is determined in this case by the medium acceleration, which is a periodic function of time, the frequency shift is also a periodic function of time, and its amplitude is determined by the amplitude of medium acceleration. The medium velocity and acceleration vectors can always be decomposed into two components, a parallel and perpendicular acoustic ray, respectively; therefore it is sufficient to estimate the longitudinal and transverse effects separately.

As to propagation in the medium of particles that perform oscillatory motion (it is assumed that α = const, see Fig. 1), the probing waves will pass through regions of the medium with different magnitude and sign of the medium velocity and acceleration. Therefore, the frequency shift does not always accumulate in time, as was in the case considered above: the maximum frequency shift will occur after a time, during which the sign of the acceleration for the longitudinal component at the receiver is not changed. This time equals half the period of the external waves. For the transverse effect this time equals one-fourth period, since the direction of the transverse motion of the medium (to one side or the opposite one) does not affect the sign of the frequency shift. In other words, to estimate the magnitude of the effect the length L must be chosen as $L = \lambda/2 = \pi v_0/\Omega$ for the longitudinal component and $L = \lambda/4 = \pi v_0/2\Omega$ for the transverse component of the frequency shift.

As is well known (see, e.g., [9]), the medium velocity and acceleration in a sound wave are related with the sound pressure by the relations

$$u = u_0 \cos \Omega\tau = \frac{\Delta p_0}{\rho_0 v_0} \cos \Omega\tau, \tag{26}$$

$$\frac{du}{d\tau} = -\frac{\Delta p_0}{\rho_0 v_0} \Omega \sin \Omega\tau. \tag{27}$$

Taking into account what was mentioned above, relations (26), (27), and using the equations of [4] for the longitudinal and transverse effects, as well as the condition $u_0 \ll v_0$, we obtain (the minus sign is omitted)

$$\Delta f_{\parallel} = f_0 \frac{\pi \Delta p_0}{\rho_0 v_0^2} \sin \Omega\tau, \tag{28}$$

$$\Delta f_{\perp} = f_0 \frac{\pi}{4 v_0^2} \left(\frac{\Delta p_0}{\rho_0 v_0} \right)^2 \sin 2\Omega\tau. \tag{29}$$

It is seen from (28), (29) that oscillatory motion of the medium generates a frequency-phase modulation of the probing waves, while the frequency shift is a periodic function with the frequency of the external wave for the longitudinal component of the oscillatory motion of the medium, and twice the frequency for the transverse component. Using for water the data given above, we obtain the following values for the maximum frequency shifts:

$$(\Delta f_{\parallel})_{max} \approx 1.4 \cdot 10^{-3} \text{ Hz}; \ (\Delta f_{\perp})_{max} \approx 1.5 \cdot 10^{-13} \text{ Hz}.$$

For air the convective effect is significantly larger. For the condition chosen for air we obtain

$$(\Delta f_{\parallel})_{max} \approx 2.2 \text{ Hz}; \ (\Delta f_{\perp})_{max} \approx 4 \cdot 10^{-5} \text{ Hz}.$$

Thus, the estimates made provide an idea on the nature and extent of deviation from the superposition principle for various media, and also show the possibility of experimental determination of these deviations.

NOTATION

v, velocity of acoustic waves in the medium; v_0, same quantity in a medium unperturbed by an external acoustic field; ρ and β, density and compressibility of the medium; f_0, probing wave frequency; F and λ, frequency and wavelength of the external waves; x, coordinate; τ, time; φ, initial phase; L, acoustic base; τ_d, delay time; p, pressure; Δp_0 and ΔT_0, amplitudes of pressure and temperature increments in the acoustic wave; α_T, bulk expansion coefficient; c_p, specific heat at constant pressure; and T, temperature.

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ELECTRICAL AND THERMAL STRUCTURE OF THE ARGON ARC OF A TWO-JET PLASMATRON

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The thermal structure and the distribution of the axial electric field strength of the arc of a two-jet plasmatron were investigated experimentally.

Characteristic features of two-jet plasmatrons (plasma generators) are the high thermal efficiency, up to 90% [1], and the presence of an exterior stabilized electric arc. Because of this plasmatrons of such type are of great promise for application in chemical technology, heat engineering, deposition techniques, spectroscopy, etc. [2-4]. The literature, however, contains no information about the results of investigations of two-jet plasmatrons with the heads opposite one another.

The aim of the present work was to determine the electrical and thermal structure of the arc burning in the open atmosphere between the opposite heads of a two-jet plasmatron and stabilized by argon streams at atmospheric pressure.

During the experiment the plasmatron, which was equipped with a rod-type thermionic tungsten cathode and an everlasting end-type copper anode [5], operated in the following conditions: arc current I = 100-200 A, voltage U = 100-300 V, flow rate of plasma-forming gas through each of the heads $G_a = G_c = (0.25\text{-}0.5)\cdot 10^{-3}$ kg/sec. The diameters of the head nozzles were $(5\text{-}7)\cdot 10^{-3}$ m, and the distance between the nozzle exits was varied in the range L = $(5\text{-}25)\cdot 10^{-2}$ m. These ranges of controlling parameters were optimal, and the alteration of even one of them in either direction led to destruction of the heads or to arrest of the arc.

The electric field strength along the arc axis was investigated with the aid of movable tungsten probes, which moved perpendicular to the arc axis at a speed W = 1 m/sec. The results of the measurements are presented in Fig. 1 (curves 1), which shows that the electric field was strongest near the nozzle exits and decreased with increasing distance from them: At a distance of about $4\cdot 10^{-2}$ m the field strength was $E = (6\text{-}7)\cdot 10^2$ V/m, which is a typical value for freely burning unstabilized argon arcs [6]. We can divide the arc lengthwise into five regions. The initial regions of the cathode and anode jets (in Fig. 1 from the nozzle exits to sections A and D, respectively) are characterized by an electric field strength that depends on the nozzle diameters and the argon flow rate. In these regions the arc is compressed and is stabilized by the rigid jets of plasma-forming gas issuing from the nozzles.

In the regions with fully developed flow (in Fig. 1 from section A to section B and from C to D) the field strength along the axis is constant and is practically independent of the arc burning conditions, which in view of the measured value $E = (6\text{-}7)\cdot 10^2$ V/m indicates laminar flow of the jets [7, 8]. In these regions the arc is spatially stable. The arc bends readily under the action of a transverse stream of gas, but when the stream is removed the arc resumes its former position. The length of these regions depends on the flow rate of the plasma-forming gas, the distance between the nozzle exits, and the lengths of the regions in which the jet flow is laminar.

The dimensions of the region where the anode and cathode plasma jets meet (between sections B and C in Fig. 1) depend on the distance L and the argon flow rate. Beginning at L = $15\cdot 10^{-2}$ m or more, the gas discharge in this region is unstable, and the scale of the instability increases with increase in the size of this region. The appearance and development of instability cause an increase in the voltage fluctuations on the arc, which can reach 40%

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