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Waveforms of acoustic pulses generated in a solid by a spark discharge

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The waveforms of longitudinal and Rayleigh wave pulses generated in a glass sample by a spark discharge are measured. A qualitative interpretation is given for the observed functional relationships.

The generation of sound in solids by spark breakdown of the gas layer adjacent to their surface is of enormous practical interest. This method has several indisputable advantages over other noncontacting methods of sound generation in solids (see, e.g., Refs. 5 and 6); in particular, it is easy to implement and highly efficient, and the spark source can be placed in relatively inaccessible locations of the tested structure.

One important aspect of the spark generation of sound in solids, which has received little attention to date, concerns the waveforms and characteristic durations of the generated acoustic pulses. The knowledge of these parameters is of considerable value for applications of the spark method in acoustic spectroscopy, nondestructive testing, and scientific research.

Here we give the results of quantitative measurements of the waveforms and durations of acoustic pulses generated by spark discharge and discuss the evolution of the waveforms and durations as the distance of the spark gap from the surface of the solid is increased.

The details of the measurement apparatus are described in Ref. 2. The spark source consisted of a simple relaxation oscillator circuit, in which the discharge capacitor \( C = 24,000 \, \text{pF} \) was charged by a high-voltage source \( U = 3 \, \text{kV} \), and then electrical breakdown took place in a spark gap of width \( d = 0.8 \, \text{mm} \). Breakdown occurred on an auxiliary grounded electrode. The acoustic pulses were generated in a glass sample of dimensions \( 22 \times 80 \times 180 \, \text{mm}^3 \).

A thick, longitudinally poled piezoceramic transducer of thickness 14 mm and diameter 9 mm was used as a wideband acoustic receiver; its back surface was damped by a thick copper plate. This type of receiver is known to be sensitive to wave displacements. However, the transducer operated in series with a specially designed broadband amplifier whose gain diminished smoothly from the center frequency of 15 MHz to zero, so that the transducer-amplifier system differentiated the displacement pulses in the frequency band from zero to approximately 10 MHz, where the waveform of the recorded electrical pulse actually coincided with that of the pressure or strain pulse. Our wideband amplifier with a gain of the order of 30 dB was constructed from high-power field-effect transistors and had a large dynamic range; these features are very important in the presence of the strong electrical noise created by discharge. The signal was observed on the screen of a storage oscilloscope. Pulses of longitudinal bulk acoustic waves generated in the direction of the


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normal to the surface of the sample and Rayleigh surface wave pulses were investigated. The receiving transducer was placed on the end of the sample in order to record the latter pulses.

Figure 1 shows photographs of the generated longitudinal bulk wave pulses for various distances from the spark gap to the sample surface. The distance between the free surface of the sample and the receiving transducer was 80 mm. We see that the pulses have a positive peak and a negative peak in every case, where the duration of the positive peak is estimated to be 150-200 ns. The generated pulses are observed to decrease in amplitude and broaden somewhat as the spark is moved farther away from the sample surface. The slight increase in the wave amplitude in Fig. 1b over Fig. 1a can be attributed to a change in the breakdown conditions of the spark gap at a certain distance from the sample surface; as the spark is moved farther from the surface, the pulse amplitude begins to decay as a result of dissipative losses of the weak air shock induced by the spark. 

Figure 2 shows Rayleigh wave pulses generated by spark breakdown for various distances between the spark gap and the surface. We see that the characteristic duration of the generated Rayleigh wave pulse is also estimated to be 150-200 ns. The generated pulses are observed to decay in amplitude and to broaden somewhat as the spark is moved farther away from the sample surface.

Consequently, the frequency spectra of the generated acoustic oscillations lie in the interval from zero to 4-5 MHz in both of the investigated spark excitation modes: longitudinal bulk waves and Rayleigh waves; this situation is evidently dictated by the physical processes associated with breakdown of the spark gap in our relaxation-type spark generator.

We now look briefly at the interpretation of the experimental data. According to the preceding discussion, we can assume that the electrical signal recorded on the oscilloscope characterizes the elastic wave pressure exerted on the receiving piezoelectric wafer. Since the principal mechanism of spark generation of acoustic waves in a solid is the action of the spark-induced air shock on its surface, the waveform of the recorded acoustic pulses in the solid must be determined by the waveform of the elastic pressure pulse.

We know that the shape of the pressure pulse of a strong, spherically divergent shock wave in air has the form shown in Fig. 3a (see, e.g., Refs. 8 and 9) in the idealized case of instantaneous energy release in a small region of space. The characteristic width $\Delta t$ of the pulse peak is $0.05r$ in this case, where $r$ is the distance from the leading edge to the center of the discharge region. As $r$ is increased, the strong shock wave gradually changes into a weak shock wave (acoustic wave) propagating with the speed of sound in air (Fig. 3b). The width of the positive peak of this wave stabilizes at a certain level $\Delta r_p$, and its subsequent evolution obeys the laws of nonlinear acoustics with allowance for the influence of dissipation. According to the well-known property of spherically divergent acoustic waves, \[ \int p' dt = 0, \] where $p'$ is the sound pressure; the pressure pulse in Fig. 3b also contains a negative peak.
If the height of the discharge gap above the surface of the solid is large enough for the shock wave to be converted into an acoustic wave as it approaches the surface (this happens at $h = h_0 = 1$ mm of the capacitance $C$ used in the experiment), the expected frequency of the $40$ MHz pulses generated in the depth of the solid can be determined approximately in accordance with the reciprocity principle. According to the latter, it is sufficient to analyze the transmission of a plane acoustic wave arriving from the depth of the solid across the solid–air interface. It is therefore obvious that the waveform of the discharge-generated acoustic pulse recorded in the depth of the solid should repeat the pulse shape of the weak air shock as it approaches the surface in this case.

A comparison of Fig. 3b with the photographs in Fig. 1 (in making the comparison, we must allow for the fact that Fig. 3 shows the dependence on distance rather than on time) indicates that the waveforms of the observed pulses agree qualitatively for heights $h$ in the interval from 1 to 2 mm (see Figs. 1b–1d). The existing discrepancies can evidently be attributed to the finite time of energy release in spark breakdown$^5$ in the experimental work and to the finite bandwidth of the amplifier, which no longer functions as a differentiating amplifier at spectral frequencies above about $40$ MHz (this fact can be explained, in particular, by the finite width of the leading edge of the pulse). The width of the positive peaks of the pulses in Figs. 1b–1d is practically constant, in correspondence with the previous discussion. The pulse shown in Fig. 1a ($h = 0.5$ mm), on the other hand, is characterized by a somewhat shorter duration, because it corresponds to a strong shock wave with as yet unestablished values of $\Delta \tau$. We note that the reciprocity principle is not valid for such a wave, so that caution must be exercised here in referring to quantitative comparison of the pulse waveforms in air and in the solid.

This explanation of the evolution of the waveforms of longitudinal-mode acoustic pulses can to a certain extent be applied to pulses of Rayleigh waves generated by spark discharge, although the physical pattern of the phenomenon is more complicated in the latter case. In particular, this is manifested in the onset of negative precursors in Figs. 2a–2d and in a monotonic amplitude decay of the generated pulses. Without going into a detailed discussion of these differences, we call attention to the fact that the durations of the positive peaks of the Rayleigh pulses for $h = 1$ mm and $1.5$ mm are approximately the same as the corresponding durations for longitudinal wave pulses. This is certainly natural, because the indicated durations are determined by the duration of the positive peak of the weak air shock $\Delta t_A = \Delta t_R/c$, where $c = 340$ m/s is the sound velocity in air. Inasmuch as $\Delta t_A = \Delta t_R h_0 = 0.05 h_0$, if we assume that $h_0 = 1$ mm in the experimental work, we obtain the estimate $\Delta t_A = 0.15 \mu s$, which is in good agreement with the experimentally observed durations of the positive peaks of the generated acoustic pulses. Theoretical estimates obtained with allowance for the variation of the shock wave velocity$^4$ for $h < h_0$, i.e., in the case of a strong shock wave incident on the surface, greatly underestimate the peak durations $\Delta t$ in comparison with the experimentally observed values for the generated acoustic pulses (Fig. 1a, 2a, and 2b).

This inconsistency can be attributed to many causes. First, there is the above-mentioned finite duration of the electric discharge and the finite width of the discharge region in comparison with the distance $h$. As a result, the pulse shape of a real spherically divergent air shock can differ appreciably from the ideal waveform (Fig. 2a). Moreover, certain corrections can be introduced by the specific nature of the transmission of a strong shock wave across the air–solid interface.

We have thus shown that acoustic pulses generated in a solid by means of a noncontacting spark generator operating in the self-excited regime have a positive peak and a negative peak with characteristic durations of the order of 150–200 ns. The width of the frequency spectrum of the generated pulses occupies the band from zero to 4–5 MHz with a maximum at 2–2.5 MHz. These pulse parameters are well suited to many applications in acoustic spectroscopy and nondestructive testing. In principle, however, it appears that they could be improved considerably by enhancing the performance of the discharge devices, which should be capable of generating a nearly ideal spherical shock wave at the shortest possible distances.

Normal modes of a waveguide with elastic walls

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The spectrum of normal modes of a planar waveguide with elastic walls is investigated. The waveguide is filled with an ideal compressible fluid. The elastic walls execute flexural vibrations. Absorption does not take place in the fluid or in the walls. It is shown that the dispersion relation for the investigated waveguide can have complex roots, along with purely imaginary roots of multiplicity 2 and 3. Numerical results and approximation equations for the multiple roots are given. An expression for the field of a point source in the waveguide is derived in the form of an expansion in normal modes. It is shown that the modes corresponding to purely imaginary and complex roots of the dispersion relation do not transfer energy through the cross section of the waveguide.

The normal modes of a planar waveguide with perfectly rigid or perfectly compliant walls have been studied, for example, in Ref. 1. However, if the walls of the waveguide are flat elastic plates, as in Refs. 2 and 3, the formulation of the normal modes has a number of distinctive features, which are discussed below.

Let a planar waveguide $-\infty < x < +\infty$, $-H < y < H$, be filled with an ideal compressible fluid. The sound pressure field $P(x, y)$ in the fluid is created by a source in physically real problems.

We assume a point source, so that the normal modes of the waveguide obey the inhomogeneous Helmholtz equation

$$\frac{\partial^2 P(x, y)}{\partial x^2} + \frac{\partial^2 P(x, y)}{\partial y^2} + k^2 P(x, y) = \delta(x, y-b),$$

(1)

where $k$ is the wave number, $k = \omega/c$, $c$ is the angular frequency, $c$ is the sound velocity in the fluid, and $b$ is the transverse coordinate of the source, $-H < b < H$. The dependence of the processes on the time $t$ is characterized by the factor $\exp(-i\omega t)$, which we omit everywhere.

The side walls of the waveguide are elastic plates. The Kirchhoff equations for flexural vibrations in contact with a fluid have the form

$$\frac{\partial^2 P(x, \pm H)}{\partial y^2} + \phi P(x, \pm H) = 0,$$

(2)

where $\kappa$ is the flexural wave number, $\kappa^2 = \kappa_x^2 + \kappa_y^2$, $\phi = \rho c^2/H$, $\rho$ is the material density of the plates, $c$ is their thickness, $D$ is the cylindrical stiffness, $\sigma_x = \rho_x c^2/D$, $\sigma_y = \rho_y c^2/D$, $\rho_x$ is the density of the fluid, and either upper signs or lower signs or both, can simultaneously be used in the ensuing equations.

Vacuum exists on the outside ($y = \pm H \pm 0$) of the plate.

Solving the problem by the method described in Ref. 4, we have

$$P(x, y) = -\frac{i}{2} \sum_{m=0}^{\infty} \left[ \frac{\text{sh}(\gamma m g^+ (\lambda, k, H))}{g^+ (\lambda, k, H)} - \frac{\text{sh}(\gamma m g^- (\lambda, k, H))}{g^- (\lambda, k, H)} \right] e^{im\omega t}.$$

(3)

The expression on the right-hand side of Eq. (3) represents the expansion of the pressure field in normal modes $P_n^*(x, y) = f_n^*(\lambda^*, k, y) e^{i\omega n t}$, $n = -2, -1, 0, 1, 2, \ldots$, $f_n^*(\lambda, k, y) = \text{sh}(\gamma y)/\gamma$, $n^2 = \lambda^2 - k^2$, and $\lambda_n^2$ denotes the roots of the equation

$$\frac{\partial P(x, k, H)}{\partial y} = \frac{\partial P(x, k, H)}{\partial y} = 0,$$

(4)

We assume in Eq. (3) that the requirement $\text{Im} \lambda > 0$ is satisfied by all $\lambda$ that would be situated in the upper half-plane of the complex variable $\lambda$ for $\text{Im} k > 0$, while $k$ itself is assumed to be a positive real number. On the complex plane of $\lambda$ we examine the roots of the dispersion relation $g^+ (\lambda, k, H) = 0$, which corresponds to the case of $\gamma$-symmetrical normal modes. For small values of $k$ the roots can be calculated by the approximate relations

$$\gamma_n^+ = \text{exp}((-i\pi/3)\text{sign}(\text{Im} \lambda)), n = -2, -1, 0, \gamma_n^+ = \text{sign}(\text{Re} \lambda), n = 1, 2, \ldots.$$

The numbers $\lambda_n^+$ and $\gamma_n^+$ are positive real for any positive $\lambda$. The roots $\lambda_n^+$, $\gamma_n^+$, $\ldots$, which are imaginary for small values of $k$, shift to the positive real axis in succession as the wave number $k$ attains the values $k_n^H$, $n = 1, 2, \ldots$, which are the roots of the equation

$$n^2 = \lambda^2 - k^2.$$

(5)

As $k$ increases (at fixed $H$), the number $\lambda_n$ can form a pair of multiplicity 2 (Fig. 1) with the