Optoacoustic sources of sound[1)]

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Usp. Fiz. Nauk 135, 637–669 (December 1981)

Laser generation of acoustic and shock waves in condensed media (liquids) is considered. It is shown that, depending on the volume density of the optical energy evolved in a liquid, the process of generation is mainly due to thermal expansion, surface evaporation, explosive boiling, or optical breakdown of the liquid. It is shown that liquids can be used as optoacoustic sources of sound. These have a number of advantages over the traditional acoustic radiators: remote action; absence of direct (in the traditional sense) contact between the body of a source and the medium which receives sound; ability to vary easily the geometric dimensions of an optoacoustic source and the frequency range; feasibility of constructing sound sources moving in a medium at an arbitrary subsonic, sonic, or supersonic velocity with complete elimination of flow of a medium around the body of a radiator. Optical generation of sound is in practice possible throughout the frequency range from infrasound to hypersound.

PACS numbers: 78.20.Hp, 43.35.Sx

CONTENTS

1. INTRODUCTION

Interaction of optical radiation with matter perturbs a material medium and this is accompanied by the emission of sound. There are various mechanisms of generation of sound and they depend primarily on the volume density of the energy evolved in a medium and on the conditions of evolution of this energy. These mechanisms include thermal expansion, electrostriction, surface evaporation, explosive boiling, and optical breakdown.

The thermal mechanism of generation of sound plays the dominant role in optically absorbing media at low energy densities. There is then no change in the aggregate state of a medium in the region of absorption of light and sound is generated because of expansion of the optically heated parts of the medium.

An increase in the energy density evolved in a medium enhances the importance of nonlinear effects due to an increase in the velocity of expansion of the heated part of the medium and also due to a change in the thermodynamic parameters of the medium in the process of interaction with laser radiation. More complex processes of generation of sound involving phase transitions and optical breakdown then become active.

The optoacoustic effect manifested by pressure vibrations of a gas in an enclosed chamber when illuminated with a modulated infrared radiation flux, was first reported by Bell in 1880 (Refs. 1–3). Beginning from the nineteen-forties, the effect has been used extensively in qualitative and quantitative analysis of gas mixtures and in detection of radiant fluxes,[4] and subsequently also in photoacoustic spectroscopy of solids and liquids.[5] However, because of the low efficiency of conversion of the optical energy into sound, the optoacoustic effect has become important in generation of sound only since the appearance of lasers. Prokhorov *et al.* were the first to observe shock waves resulting from the interaction between a laser beam and water.[6] Attempts have been made later to account for some

[1)] Invited paper presented at Tenth International Acoustic Congress, Sydney, Australia, 1980.

 0038-5670/81/120977-19$01.80

features of the optical generation of sound. This work is discussed in a review by Bunkin and Komissarov.[7]

Currently investigations of the optical generation of sound are proceeding on a wide front. A considerable number of papers on laser generation of sound has already been published. It is now possible to speak of optical (usually called optoacoustic) sources of sound.

Optoacoustic sources have a number of advantages over traditional acoustic radiators: remote action; absence of direct (in the traditional sense) contact with a medium receiving sound; ability to alter readily the geometric parameters of an optoacoustic antenna and the range of emitted frequencies; feasibility of constructing sound sources moving at practically arbitrary subsonic, sonic, or supersonic velocities with complete elimination of flow of a medium around the body of a radiator; ability to generate sound optically in a wide range of frequencies from very low right up to hypersonic.

2. THERMAL MECHANISM. EQUATION FOR OPTICAL GENERATION OF SOUND AND METHOD OF SOLVING BOUNDARY-VALUE PROBLEMS

We shall consider the specific case of optical generation of sound in liquids. The main features of this effect at light intensities such that the density of the optical energy evolved in a medium is small compared with the heat of evaporation can be described satisfactorily by a linear theory.

In fact, at low densities of the energy evolved in a liquid the rate of expansion of a heated region is low compared with the velocity of sound and one can use linearized equations of hydrodynamics. Moreover, in most cases of practical interest it is possible to ignore the influence of the thermal conductivity on the process of sound generation (we then have $l \gg \chi/c$, where l is the characteristic size of the heated region of a liquid, χ is the thermal diffusivity, and c is the velocity of sound in the selected liquid). It is usual also to ignore the attenuation of sound, which can always be allowed for separately as is usually done in acoustics.

We shall begin with a system of equations of hydrodynamics and of electromagnetic fields. Linearizing the system of equations and ignoring the viscosity and heat conduction effects, we obtain an inhomogeneous wave equation for the optical generation of sound[8]:

$$\Delta p-\frac{1}{c^2}\frac{\partial^2 p}{\partial t^2}=-\frac{\varkappa}{C_p}\frac{\partial Q}{\partial t}, \tag{1}$$

where p is the acoustic pressure; $\varkappa$ is the volume thermal expansion coefficient; C_p is the specific heat of the investigated liquid; c is the velocity of sound in the liquid; Q is the power density of thermal sound sources due to the absorption of the optical energy in the liquid given by $Q = \operatorname{div} \mathbf{F}$; $\mathbf{F}$ is the density of the optical energy flux in the liquid (Poynting vector).

The present author[9] suggested an effective method for solving the boundary-value problems in optical generation of sound. Solution of a boundary-value problem with the aid of the integral reciprocity relationship[10] can be reduced directly to quadratures and represented in the form

$$p(r,t)=-\frac{\varkappa}{C_p}\int\limits_{t'}\int\limits_{\Omega'}\frac{\partial Q(r',t')}{\partial t'}\tilde{p}(r',t'\mid r,t)\,d\Omega'\,dt', \tag{2}$$

provided we know the solution $\tilde{p}(r',t'\mid r,t)$ of an auxiliary self-conjugate diffraction problem of the field of a point source when the source is located at that point $r(x,y,z)$ in a medium where the field of thermal sources of sound has to be determined. Here, Ω is the region of action of thermal sources of sound.

In most cases the interest is centered on the field in the Fraunhofer zone. It is then sufficient to know the solution of the auxiliary problem of diffraction of a plane acoustic wave. This solution is in most cases either already known or it can be found quite readily.

3. GENERATION OF MONOCHROMATIC SOUND

We shall assume that a laser beam traveling from the upper half-space in the positive direction of the z axis is incident on the free surface of a liquid and the equation describing this surface is $z=\xi(x,y)$, whereas the intensity of light varies harmonically at a frequency ω (Fig. 1). Then, thermal sources of sound form in the liquid and the power density of these sources is given by

$$Q(x,y,t)=A\mu J(x,y)\exp\{-\mu[z-\xi(x,y)]\}(1+m\cos\omega t),$$

where $J(x,y)$ is the distribution of the intensity of light in the laser beam which is usually assumed to be Gaussian $J(x,y)=J_0\exp-(x^2+y^2)/a^2$; m is the modulation index defined by $0\leqslant m\leqslant 1$; μ is the optical absorption coefficient of the liquid; A is the optical transmission coefficient of the surface of the liquid; a is the radius of the illuminated spot on the surface of the liquid.

We shall now consider the characteristics of an acoustic field of an optoacoustic source in the Fraunhofer zone in some special cases.

A. Half-space with a plane boundary $z=\xi(x,y)\equiv 0$

The auxiliary solution is of the form

$$p(r',r)=\frac{\exp(ikr)}{4\pi r}\{\exp[-i(\alpha x'+\beta y'+\gamma z')]-\exp[-i(\alpha x'+\beta y'-\gamma z')]\},$$

where $\alpha^2+\beta^2+\gamma^2=k^2$; $r^2=x^2+y^2+z^2$; $k=\omega/c$ is the wave number of sound in liquid; the factor $\exp(-i\omega t)$ is, as usual, omitted throughout our treatment.

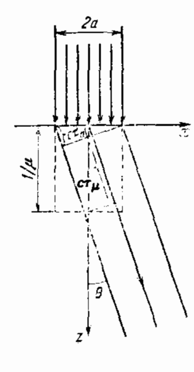

FIG. 1. Geometry of the problem.

Equation (2) yields the acoustic pressure[11,12]

$$p(r) = -\frac{\omega m \varkappa A J_0 a^2}{2C_p} \frac{\exp(ikr)}{r} \frac{\mu k \cos\theta}{\mu^2 + k^2\cos^2\theta} \exp\left(-\frac{k^2a^2}{4}\sin^2\theta\right); \quad (3)$$

here θ is the angle between the direction of incidence of the laser beam and the direction from the point of observation to the origin of the coordinate system; r is the distance from the point of observation to the origin.

We can see from Eq. (3) that the amplitude of the acoustic pressure increases on increase in the laser power proportionally to J_0a^2 and it also increases on increase in the frequency and modulation index. Directionality of the acoustic emission depends on the parameters ka and $k\mu^{-1}$.

If $k\mu^{-1} \leqslant 1$ and $ka \ll 1$, dipole emission of sound is observed because under these conditions a monopole source acts on the free surface of a liquid and the radiation field of such a source is of the dipole type due to the influence of the free surface Fig. 2(a) .

If $k\mu^{-1} \gg 1$ and $ka \ll 1$, sound is emitted mainly along the surface. A set of volume sources forms a thin (in the transverse direction) vertical antenna which is long compared with the acoustic wavelength and which is directed along the laser beam Fig. 2(b) .

If $k\mu^{-1} \ll 1$ and $ka \gg 1$, an acoustic antenna is in the form of a disk whose diameter is much greater than the acoustic wavelength. The radiation is directed mainly along the laser beam Fig. 2(c) .

An analysis of Eq. (3) also shows that the optimal conditions for the emission of sound as a result of laser excitation are obtained[11] when $k = \mu$. This imposes certain conditions on the laser frequency (wavelength).

B. Half-space with a ruffled uneven surface $z = \xi(x,y)$

Under real conditions the surface of a liquid is frequently ruffled (uneven) and this may be due to a number of reasons. We shall assume that the unevenness is large compared with the acoustic wavelength, i.e., we shall assume that $k\sigma \gg 1$, where

$$\sigma = \sqrt{\langle|\xi(x, y)|^2\rangle}, \quad \langle\xi(x, y)\rangle = 0,$$

and $\langle \ldots \rangle$ denotes as usual the statistical averaging operation.

Using an approximation similar to the Kirchhoff approximation in the theory of diffraction of waves by an uneven surface, we can write down the solution of the auxiliary diffraction problem

$$p(r', r) = -i\frac{\exp(ikr)}{2\pi r} \exp\{-i[(\alpha x' + \beta y' - \gamma\xi(x', y')]\} \times \sin\{\gamma[z' - \xi(x', y')]\} + \psi',$$

where ψ' is some unimportant correction.

If displacements of the surface are statistically homogeneous, isotropic, and obey the normal distribution law, we find that the average pressure is given by[13,14]

$$\langle p(r)\rangle = -\frac{m\varkappa\omega A}{2C_p} J_0a^2 \frac{\exp(ikr)}{r} \frac{\mu k\cos\theta}{\mu^2 + k^2\cos^2\theta} \times \exp\left(-\frac{k^2a^2}{4}\sin^2\theta\right)\exp\left(-\frac{q^2}{8}\right), \quad (4)$$

where $q = 2k\sigma\cos\theta$ is the Rayleigh parameter of a ruffled uneven surface.

The average field is directional. For angles of observation $\theta < \theta_0 = \tan^{-1}(\sqrt{2}\,\sigma/a)$ the directionality is affected considerably by the unevenness of the surface (Rayleigh parameter), whereas in the range $\theta > \theta_0$ the directionality depends mainly on the ratio of the acoustic wavelength to the dimensions of an optical spot, and the unevenness of the surface plays practically no role in the formation of the field of an optoacoustic source.

The average intensity of the field in the case when the displacements of the surface obey a two-dimensional normal distribution law is[14]

$$\langle|p(r)|^2\rangle = \left(\frac{mc\varkappa A}{2C_pr}\right)^2 (J_0a^2)^2 \frac{\mu^2k^4\cos^2\theta}{(\mu^2 + k^2\cos^2\theta)^2} \frac{\exp\left[-\left(\frac{\rho_0}{\sigma}\right)^2 \frac{\Delta^2\sin^2\theta}{4(\Delta^2\cos^2\theta + 1)}\right]}{1 + \Delta^2\cos^2\theta}, \quad (5)$$

where $\Delta = \sqrt{2}\,ka\sigma/\rho_0$, and ρ_0 is the radius of the spatial correlation of displacements of a ruffled surface.

We shall now consider special cases.

Let us first assume that $\Delta \ll 1$. Then, since $k\sigma \gg 1$ (strongly ruffled surface), the condition $\Delta \ll 1$ corresponds to the optical spot radius being small compared with the correlation radius of the surface unevenness, i.e., $a \ll \rho_0$, and in the limit $\Delta \to 1$ the average intensity approaches that of an acoustic field resulting from the absorption of modulated laser radiation in a liquid with an unruffled plane surface. In other words, if $a \ll \rho_0$, the influence of the unevenness on the average intensity of the field can be ignored.

If the radius of an optical spot is comparable or greater than the correlation radius of the surface unevenness $a \geqslant \rho_0$, i.e., if $\Delta \gg 1$, then in the $k\sigma \gg 1$ case we have

$$\langle|p(r)|^2\rangle \approx \left(\frac{mc\varkappa A}{2C_pr}\right)^2 (J_0a^2)^2 \frac{k^4}{a^2(\mu^2 + k^2\cos^2\theta)^2} \left(\frac{\mu\rho_0}{k\sigma}\right)^2 \times \exp\left[-\frac{1}{4}\left(\frac{\rho_0}{\sigma}\right)^2 \mathrm{tg}^2\theta\right]. \quad (6)$$

In this case the directionality is governed by the scale of the surface unevenness. The intensity of sound decreases by a factor e when the angle of observation is $\theta = \tan^{-1}(\sqrt{2}\,\sigma/\rho_0)$.

These results can be illustrated graphically. Figure 3 shows the angular dependence of the average pressure for an optoacoustic source with the wave dimensions $ka = \sqrt{10}$. Curves 1–4 correspond to the following val-

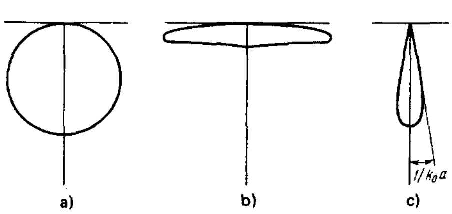

FIG. 2. Angular distributions of the radiation emitted from an optoacoustic source in certain limiting cases.

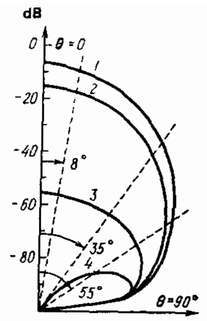

FIG. 3. Polar diagram of the average pressure.

ues of the parameter σ/a: 0.01, 0.1, 0.5, 1. The dashed lines give the angles $\theta=\theta_0=\tan^{-1}(\sqrt{2}\,\sigma/\rho_0)$. Figure 4 shows the characteristics of the average field intensity calculated for the same source. Curves 1–3 correspond to the following values of Δ^2: 0.1, 1, 10. The dashed line represents the limit of the angular width of the directionality characteristic $\theta_0=\tan^{-1}(\sqrt{2}\,\sigma/\rho_0)=14^\circ$ for the selected ratio σ/ρ_0.

The influence of a small (compared with the acoustic wavelength) unevenness of the surface of a liquid on the optical generation of sound is considered in Ref. 15 and the influence of a boundary in the form of a superposition of large and small unevennesses is analyzed in Ref. 16. In the latter case (known as the two-scale model of an uneven surface) it has been established that the expression describing the average field is

$$\langle p(r)\rangle = p_0(r)\,f(-\gamma)\,f_1(\theta),$$

where $f(-\gamma)$ is the characteristic function of a random quantity describing the large-scale unevenness, whereas $f_1(\theta)$ is a function describing the influence of small-scale unevenness which depends on the nature of the energy spectrum of this unevenness. An explicit expression for $f_1(\theta)$ is obtained in Ref. 16. It should be noted that $p_0(r)$ describes an acoustic field in the half-space occupied by a liquid with a plane surface. The above expression is obtained on the assumption that large and small unevennesses of the surface are statistically independent. In particular, when large-scale random displacements of the surface of a liquid obey the normal distribution law, it is found that $f(-\gamma)=\exp(-q^2/8)$, where, as before, q is the Rayleigh parameter.

It is interesting to consider the question of the efficiency of conversion of the optical energy into sound. Following Ref. 11, this efficiency can be described by the ratio of the total power of sound to the power of laser radiation. For example, in the case of a wide laser beam ($ka \gg 1$), it follows from Eq. (3) that the efficiency is

$$\eta \approx \frac{c}{\rho}\left(\frac{Am\varkappa}{2C_p}\,\frac{\mu k}{\mu^2+k^2}\right)^2 J_0.$$

The maximum efficiency corresponds to $\mu \approx k$ and it is given by

$$\eta_{\max} \approx \frac{c}{\rho}\left(\frac{Am\varkappa}{4C_p}\right)^2 J_0.$$

It thus follows that the efficiency of conversion is directly proportional to the intensity of optical radiation. It should be noted that for water we have $\eta_{\max}\approx 5\cdot 10^{-12} J_0$, where J_0 is in watts per square centimeter.

An increase in the intensity of light can increase greatly the efficiency of conversion of the optical energy into sound. However, this is true only as long as the thermal mechanism of the excitation of sound is active, i.e., as long as the density of the optical energy evolved in a liquid is low compared with the heat of evaporation.

C. Generation of monochromatic sound in an inhomogeneous liquid

We have considered above the main laws governing the optical generation of sound in a homogeneous liquid. However, real liquids are to a greater or smaller extent inhomogeneous. Inhomogeneities may sometimes have a considerable influence either directly on the process of optical generation of sound or on the characteristics of an acoustic field in the course of propagation of acoustic waves in a medium, or there may be simultaneous manifestation of both effects of the inhomogeneous properties of a liquid.

Guided propagation of acoustic waves may be important. Some features of the optical generation of sound in a liquid waveguide are considered in Ref. 17. It is shown that, depending on the wavelength of optical radiation (value of the coefficient μ) and laser beam radius, one can expect optimal excitation of some specific normal wave traveling in a waveguide.

Characteristics of the excitation of monochromatic sound in the case of absorption of modulated laser radiation in a liquid half-space with a homogeneous layer of a different liquid on the surface are considered in Ref. 18. The presence of a layer of another liquid or of the

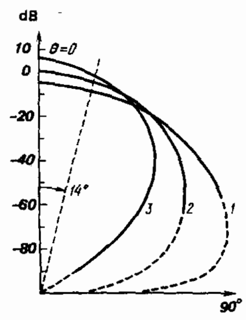

FIG. 4. Polar diagram of the average intensity of the acoustic field.

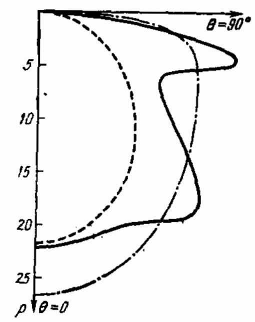

FIG. 5. Augular distribution of the radiation emitted from an optoacoustic source in water with a surface layer of benzene.

same liquid but heated or containing gas bubbles or small inclusions may alter greatly the intensity and angular distribution of the acoustic field. For example, Fig. 5 shows the angular distributions of sound emitted by an optoacoustic source corresponding to the case when a layer of benzene is present on the surface of water. It is assumed that the coefficient of absorption of light in water is $\mu = 0.18\ \mathrm{cm}^{-1}$, whereas the corresponding coefficient of benzene is $\mu = 2.3\ \mathrm{cm}^{-1}$; the ratio $\varkappa/C_p$ for water is 4.7×10^{-12} g/erg, whereas for benzene it is 62.3×10^{-12} g/erg; the ratio of the densities of benzene and water is 0.88 and the velocities of sound in water and benzene are in the ratio 0.89. The radius of an optical spot is 0.4 cm. The curves are calculated for three values of the benzene layer thickness. The dashed curve corresponds to a benzene layer 0.33 cm thick, the chain curve corresponds to 0.66 cm, and the continuous curve to 1.99 cm. The pressure amplitude is normalized to $A\omega m J_0 a^2/2r\ \mathrm{g\cdot cm^{-1}} \times \mathrm{sec}^{-2}$. The angular distribution of an optoacoustic source of sound consisting of water without a layer of benzene is, under the same conditions, strongly elongated along the surface and the amplitude of the acoustic pressure in the direction of propagation of the laser beam ($\theta = 0$) is approximately two orders of magnitude lower than in the cases when a layer is present.

Generation of sound by modulated laser radiation in a liquid in contact with a thin elastic plate is investigated theoretically in Ref. 19. Two limiting cases are considered: (a) an optically transparent plate so that laser radiation is absorbed in the liquid; (b) laser radiation is absorbed in the surface layer of the plate and the acoustic field in the liquid is generated by the vibrations of the plate. In both cases the vibrations of the plate exert the most important influence on the characteristics of the acoustic field in the liquid along directions which are governed by the conditions of "spatial resonance" of longitudinal or flexural vibrations of the plate. By way of example, Fig. 6 gives the angular distributions of the acoustic pressure in water with a free surface (curve 1) and when the water surface is in contact with a thin steel plate (curve 2). The calculations are carried out for the following parameters: plate thickness 1 mm, laser spot radius 1 mm, modulation frequency 10^6 Hz, laser power 1 W.

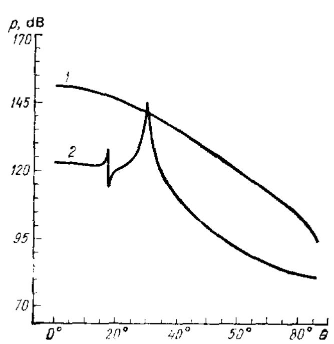

FIG. 6. Angular distribution of the radiation emitted by an optoacoustic source in water: 1) free surface; 2) thin steel plate on the surface of water.

The maxima in the angular distribution of the radiation in the case when there is a plate on the surface of water (curve 2) are due to spatial resonances of longitudinal ($\theta = 16°$) and flexural ($\theta = 30°$) vibrations in the plate.

4. EXCITATION OF SOUND BY LASER PULSES

A. Generation of sound by a single laser pulse

Excitation of sound by laser pulses is of special interest because pulsed lasers can provide high-power radiation. The process of thermooptic generation of sound by laser pulses has been studied extensively both theoretically and experimentally. The first theoretical treatments[20,21] have been concerned with one-dimensional problems and therefore it is not surprising that in some cases the experimentally observed acoustic pulses have been of very different shape from those predicted theoretically. This discrepancy is attributed in Ref. 22 to nonlinear effects in the generation and propagation of acoustic pulses in liquids. Attempts have been made to obtain numerical solutions of nonlinear equations of hydrodynamics.[23] However, detailed theoretical studies[24,25] have shown that at low optical energy densities in a medium the linear theory describes correctly the actual process of generation of sound by laser pulses and the results obtained in such theoretical treatments are in good agreement with the results of earlier experiments, including those reported in Ref. 22, and also with the results of special tests.[26-29]

We shall now consider in greater detail the main laws governing generation of sound by laser pulses in liquids. If we use the Fourier transformation and the solution for monochromatic acoustic vibrations discussed above, we obtain the following expression for the acoustic pressure in the far-field zone in the case when a liquid is perturbed by a laser pulse[25]

$$p(r, t) = -\frac{A\varkappa J_0 a^2}{4\pi C_p \tau_\mu^2 r}\left[\int_{-\infty}^{+\infty} \exp\left(-\frac{u^2 s^2}{4} + iuv\right) F\left(\frac{u}{\tau_\mu}\right) \mathrm{d}u - \int_{-\infty}^{+\infty} \exp\left(-\frac{u^2 s^2}{4} + iuv\right) \frac{F(u/\tau_\mu)}{1+u^2}\, \mathrm{d}u\right], \qquad (7)$$

where

$$u = \omega\tau_\mu, \quad v = \left(\frac{r}{c} - t\right)\tau_\mu^{-1}, \quad s = \frac{\tau_a}{\tau_\mu} \text{ and } \tau_\mu = \frac{\cos\theta}{\mu c}, \quad \tau_a = \frac{a\sin\theta}{c}$$

are the characteristic delay times of sound from elementary sources in the vertical and horizontal cross sections of the region occupied by thermal sources (Fig. 1), and

$$F\left(\frac{u}{\tau_\mu}\right) = F(\omega) = \int_{-\infty}^{+\infty} f(t)\, e^{i\omega t}\, \mathrm{d}t$$

is the spectrum of a laser pulse.

An analysis of Eq. (7) is based on the fact that the spectral width of the functions occurring in the integrand depends on the characteristic time scales, which are the pulse duration τ, and the delay times τ_a and τ_μ. In fact, we find that for the laser pulse spectrum we have $\omega \leqslant c_1/\tau$, for an exponential function we have $\omega \leqslant c_1/\tau_a$, and for a rational function we have $\omega \leqslant c_1/\tau_\mu$, where c_1 is a constant.

It follows from Eq. (7) that in the case of long laser pulses when $\tau \gg \tau_a$ and $\tau \gg \tau_\mu$, the shape of an acoustic pulse is governed by the second derivative of the laser pulse envelope:

$$p(r,\ t) \approx -\frac{A\varkappa J_0 a^2}{4\pi C_p r}\tau_\mu f''\left(t-\frac{r}{c}\right). \tag{8}$$

If $\tau \gg \tau_a$ but $\tau \ll \tau_\mu$, i.e., when the region of effective heat evolution (i.e., the optoacoustic antenna) is in the form of a narrow cylinder, the acoustic signal is a rarefaction pulse whose shape repeats the envelope of an inverted laser pulse $f(t)$ with a positive correction proportional to a small parameter τ/τ_μ:

$$p(r,\ t) = -\frac{A\varkappa J_0 a^2}{4\pi C_p r \tau_\mu}\left\{f\left(t-\frac{r}{c}\right)-\pi\frac{\delta}{\tau_\mu}\exp\left[-\left|t-\frac{r}{c}\right|\right]\right\}, \tag{9}$$

and $\delta = \int_{-\infty}^{+\infty} f(t)dt$ is the area under the laser pulse.

It should be noted that in the direction of incidence of the laser beam ($\theta = 0$) the expressions (8) and (9) are valid for any radius of the optical spot ($\tau_a = 0$ for any value of a).

If $\tau \ll \tau_a$, we have

$$p(r,\ t) = -\frac{A\varkappa J_0 a^2 \delta}{8C_p r\tau_\mu^2}\left\{\frac{4}{\sqrt{\pi}s}\exp\left(-\frac{\gamma^2}{s^2}\right)-\exp\left(\frac{s^2}{4}\right)\right.$$
$$\left.\times\left[\exp(-\gamma)\mathrm{Erfc}\left(\frac{s}{2}-\frac{\gamma}{s}\right)+\mathrm{Erfc}\left(\frac{s}{2}+\frac{\gamma}{s}\right)\exp(\gamma)\right]\right\}, \tag{10}$$

where $\mathrm{Erfc}(z) = (2/\pi)\int_z^\infty \exp(-t^2)dt$ is the complementary error function; $\gamma = [t-(r/c)]/\tau_\mu$.

The shape of the acoustic pulse turns out to be universal (Fig. 7), i.e., it is independent of the envelope of the laser pulse but it is governed by the ratio s of the characteristic times. The shape of the laser pulse determines only the amplitude of the acoustic signal. This signal consists of two compression stages separated by a rarefaction stage, and the signal is symmetric relative to the moment $t-(r/c)$.

B. Excitation of sound by a train of laser pulses

Investigations of an acoustic field excited in a liquid by a train of laser pulses were reported in Refs. 30 and 31. A special feature of such generation is a large number of acoustic harmonics. The width of the excited spectrum is governed by the repetition frequency of laser pulses, their duration, and absorption coefficient of laser radiation in the selected medium. The spectrum of the acoustic field may include harmonics generated optimally, compared with others if for a given spectral component the conditions of optimal excitation are satisfied, i.e., if $k_n \approx \mu$, where k_n is the wave number of the n-th acoustic harmonic. When sound is excited by a sequence of laser pulses, the efficiency of conversion of the laser radiation into sound increases.

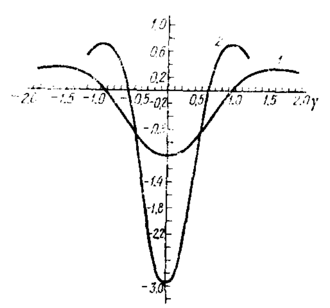

FIG. 7. Acoustic signal generated by a short laser pulse.

5. GENERATION OF SOUND IN A MOVING MEDIUM. ACOUSTIC FIELD OF A MOVING OPTOACOUSTIC SOURCE

The influence of motion of a liquid on optical generation of sound was considered in Ref. 32. Allowance was made for waves on the surface of the liquid.

Detailed studies have been made on the field of a moving optoacoustic source.[33-37] The characteristics of generation of sound in a liquid by a laser pulse of arbitrary shape in the case when a laser beam is moving along the surface of a liquid were analyzed theoretically in Ref. 37. Hardly any limitations were imposed on the velocity of the laser beam or on the shape of its trajectory. It was simply assumed that the trajectory of the beam was located on a finite region of the surface of a liquid and the acoustic field was investigated in the far-field zone relative to this region. The following cases were considered in detail: (a) rectilinear and uniform motion of the laser beam along a finite trajectory; (b) oscillatory motion of the beam; (c) uniform motion of the beam on a circle.

The results showed that the generation of sound by a uniformly and rectilinearly moving pulsed optoacoustic source occurs in the same way as in the case of a source at rest excited by a laser pulse compressed by a factor $|1-\tilde{M}|$ with an effective duration $\tau|1-M|$ and shape described by

$$f\left(\frac{t}{|1-\tilde{M}|}\right)|1-\tilde{M}|^{-1},$$

where τ is the duration of the laser pulse; $\tilde{M} = (V/c) \times \sin\theta\cos\varphi$; V is the velocity of the laser beam on the liquid surface; c is the velocity of sound in the liquid; θ and φ are the angular coordinates representing the direction of the point of observation.

It follows that practically all the conclusions and results given earlier for the generation of sound by an immobile pulsed laser beam apply also to the case of a moving beam if we consider an effective laser pulse.

If the laser beam moves at a supersonic velocity along the surface of a liquid, i.e., if $\tilde{M} > 1$, the effective laser pulse is not only compressed but also inverted with respect to the time axis relative to the real pulse. This is due to the fact that for certain directions of observation the acoustic perturbations created by supersonically traveling laser beam arrive earlier at the point of observation.

The envelope of an acoustic signal in the Cherenkov direction $(V/c)\sin\theta_1\cos\varphi_1 = 1$ observed in the case of supersonic motion of the source is independent of the shape and duration of the laser pulse. It follows from the above discussion that this is due to the fact that the acoustic field is generated by a very short "compressed" laser pulse.

However, the envelope of the laser signal in directions other than or even close to that of the Cherenkov emission is related in a specific way to the nature of the laser pulse. A characteristic situation occurs in the case of, for example, intrapulse modulation of a laser pulse of frequency ω_0. An acoustic signal in the Cherenkov direction is independent of the shape of the laser pulse and of the nature of intrapulse modulation, but the acoustic signal along directions close to the Cherenkov is a pair of pulses in the case of intrapulse laser modulation and this pair of acoustic pulses represent the response to the beginning and end of the laser pulse. Each of these acoustic pulses consists of compression and rarefaction stages and the interval between the acoustic pulses is of the order of $2|1-\tilde{M}|\tau$. If the inequality $\omega_0 < 2|1-M|/\tau$ is obeyed, the interval between the acoustic pulses is filled by an almost sinusoidal signal at the Doppler frequency $\omega_0/|1-\tilde{M}|$. By way of illustration, Fig. 8 shows the envelope of an acoustic pulse in the Cherenkov direction (curve 1) and the acoustic pulses corresponding to the beginning and end of a laser pulse (curves 2 and 3, respectively).

We shall conclude by noting that the acoustic effects similar to those described above may be observed when an optical focus moves in a nonlinear medium under conditions of self-focusing of laser radiation.[38]

6. OPTOACOUSTIC TRANSITION RADIATION

If there is a change in the properties of a medium on the surface or in the surface layer of a liquid (or solid) and a laser beam of constant intensity crosses this region, the acoustic signal which is generated can be called the transition radiation by analogy with the transition radiation well known from electrodynamics. In particular, the optoacoustic transition radiation should be observed when a laser beam crosses a boundary between two media and it may be due to difference of any of the following parameters of the two media: velocity of sound, density, thermodynamic properties, etc. This transition radiation should also appear during motion of the laser beam when an optical spot emerges from behind the screen. If light in the laser beam is intensity-modulated, the optoacoustic transition radiation which appears in such cases "distorts" the acoustic signal which is due to this modulation and which appears as a result of the absorption of light in a liquid. Some properties of the optoacoustic transition radiation were investigated in Refs. 39 and 40; attention to this effect was drawn for the first time in these investigations.

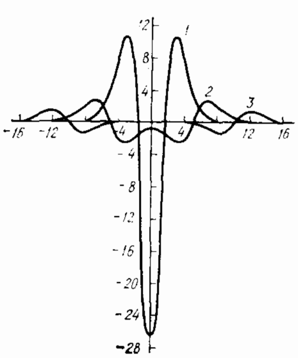

FIG. 8. Envelope of an acoustic signal in the Cherenkov direction (curve 1) and pulses at the beginning (2) and end (3) of the laser pulse in the case of supersonic motion of a pulsed optoacoustic source.

7. OPTOACOUSTIC CONCENTRATORS OF SOUND

The region of absorption of light in a medium can be shaped in a specific manner, which makes it possible to control the divergence of an acoustic beam excited in a medium. Optoacoustic concentrators in which a laser beam is absorbed in a thin layer of a liquid confined by an optically and acoustically transparent solid spherical or cylindrical shell have been proposed and investigated both theoretically and experimentally.[41,42] Clearly, optoacoustic concentrators provide an opportunity of generating strong acoustic signals at ultrasonic and hypersonic frequencies.

8. SOME EXPERIMENTAL RESULTS ON LASER GENERATION OF SOUND UNDER THE THERMAL MECHANISM CONDITIONS

A qualitative agreement between the experimental and theoretical results was already obtained in Ref. 43. Measured dependences of the efficiency of laser generation of sound on the parameter $\varkappa$ were reported in Refs. 44 and 45. The angular characteristics of optoacoustic sources were also measured in the near- and far-field zones.

The results reported in the cited papers and the experiments carried out by other authors have demonstrated that the above theoretical ideas are in good agreement with the experimental observations.

For example, it follows from the theory that the amplitude of the acoustic pressure rises linearly on increase in the optical radiation power. This is confirmed by the experimental results. The continuous line in Fig. 9 is the theoretical dependence and the points are the experimental results.[46] The ordinate gives the acoustic pressure on the axis of an optoacoustic source, i.e., in the direction of propagation of a laser beam. The pressure is normalized to a distance of 1 m and a pressure of 10^{-6} Pa. The abscissa shows changes in the optical radiation power in kilowatts.

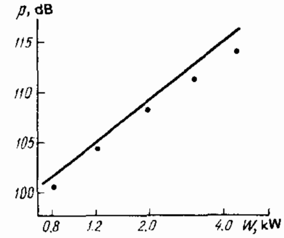

FIG. 9. Dependence of the acoustic pressure on the axis of an optoacoustic source on the laser power. The continuous line is theoretical and the points represent the experimental results from Ref. 46.

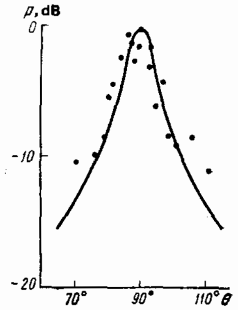

FIG. 10. Angular distribution of the acoustic field in the case of laser excitation of sound in water. The continuous curve is theoretical and the points are the experimental results taken from Ref. 46.

The experiments reported in Ref. 46 were carried out in lake water. A neodymium laser ($\lambda_l = 1.06\ \mu$) was operated in the pulsed regime and intrapulse modulation of the optical radiation intensity was applied. The modulation frequency was such that sound generated in water was quasimonochromatic.

Figure 10 shows the theoretical (continuous curve) and experimental (points) results on the angular distribution of the acoustic field generated by laser radiation.[46] In this case the frequency of sound was $f = 50 \times 10^3$ Hz, the regime was such that $ka \ll 1$, the measurements were carried out at a distance of 16.8 m, and the optical absorption coefficient of water was $\mu = 15.7\ \mathrm{m}^{-1}$.

Figure 11(a) shows acoustic pulses predicted by the theory on the assumption that a laser pulse is rectangular[24] and these are compared with the experimental results[22] in Fig. 11(b). The calculations were carried out allowing for the experimental conditions in Ref. 22.

Figure 12 shows oscillograms of an acoustic signal obtained in the far-field zone as a result of excitation with short laser pulses.[29] The acoustic pulses were excited in an aqueous solution of copper sulfate. Variation of the concentration of this solution could alter the absorption coefficient of light. The source of light was a ruby laser emitting pulses of 4×10^{-8} sec duration and 0.08–0.1 J energy. Oscillograms were obtained for the following values of the parameters: (a) $\mu = 0.3\ \mathrm{cm}^{-1}$, $a = 2.3$ cm, $\theta = 60°$; (b) $\mu = 0.8\ \mathrm{cm}^{-1}$, $a = 2.3$ cm, $\theta = 60°$; (c) $\mu = 1\ \mathrm{cm}^{-1}$, $a = 3$ cm, $\theta = 40°$. The shape of the acoustic pulses was in agreement with

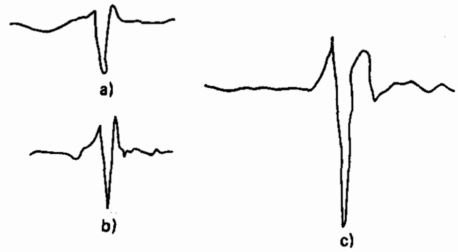

FIG. 12. Oscillograms of acoustic signals in the far-field zone, generated by laser pulses in an aqueous solution of copper sulfate.[29]

that calculated from Eq. (10) for short laser pulses (Fig. 7).

The results of the first experimental investigations of the acoustic field of a moving optoacoustic source are presented in Refs. 47 and 48. An experimental study of the optoacoustic transition radiation was reported recently.[49] The experimental results[47–49] were in agreement with the theoretical predictions. For example, Fig. 13 gives the angular distributions in the horizontal plane of a moving optoacoustic source.[48] This source was established in water by laser radiation of wavelength 1.06 μ. The optical absorption coefficient of water at this wavelength was $\mu = 0.17\ \mathrm{cm}^{-1}$. The modulation frequency of laser radiation was $f = 10^5$ Hz. A laser pulse carried an energy $E \approx 0.2$ J and its duration was $\tau = 0.5$ msec. The transverse distribution of the intensity in the laser beam on the surface of water was nearly Gaussian. The radius of the optical spot on the surface of water was $a \approx 0.25$ cm. The continuous curves in Fig. 13 are the results of calculations and the points are the experimental data. We can see that the angular distribution in the horizontal plane of a moving optoacoustic source is elongated in the direction of motion. An increase in the velocity of motion increases the amplitude of the sound emitted in the forward direction and decreases the amplitude in the opposite direction.

The motion of an optoacoustic source influences the angular distribution in the vertical plane. The distribution becomes narrower in the direction of motion of the source and broader at right-angles to this direction. Figure 14 shows the results of a determination of the angular distribution in the vertical plane in the direction of motion of an optoacoustic source of the parameters given above when the velocity of motion was $V = 0.3c$, where c is the velocity of sound. The theoretical results are represented by the continuous

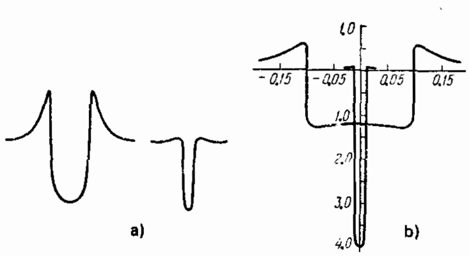

FIG. 11. Acoustic pulses recorded experimentally[22] (a) and predicted theoretically[24] (b).

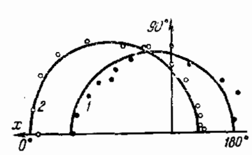

FIG. 13. Angular distribution of the radiation in the horizontal plane of a moving optoacoustic source[48]: 1) $V = 0.15c$; 2) $V = 0.47c$ (Ref. 48).

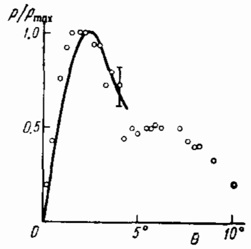

FIG. 14. Angular distribution of the radiation in the vertical plane of a moving optoacoustic source recorded in the direction of motion ($V = 0.3c$).[48]

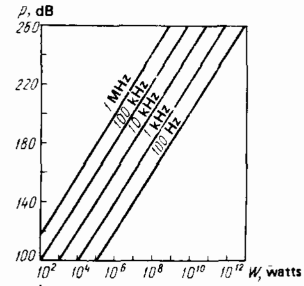

FIG. 16. Dependence of the acoustic pressure on the frequency of sound and on the power of optical radiation (neodymium laser) in pure water.[46]

curve. The dependence of the half-width of the angular distribution on the velocity is plotted in Fig. 15 (Ref. 48).

We shall conclude this section by quoting the results of calculations of the dependence of the acoustic pressure in pure water on the velocity of sound and on the optical radiation power in water of a neodymium laser (Fig. 16, which is plotted on the basis of the results reported in Ref. 46). The abscissa gives the optical power in water in watts and the ordinate gives the acoustic pressure in decibels relative to 10^{-6} Pa, reduced to 1 m in the direction of the laser beam. These results are some measure of the efficiency of conversion of the optical energy into sound in the case of the thermal mechanism.

9. OPTICAL GENERATION OF SOUND. NONLINEAR EFFECTS

A. Nonlinear thermodynamic effects

At moderate densities of the energy evolved in a medium we can expect nonlinear effects due to the temperature dependences of the thermodynamic parameters of the medium. If these effects are small, it is possible to develop an approximate theory of the optical generation of sound allowing for nonlinear thermodynamic effects.[50,51]

Such a theory is founded essentially on replacement of the right-hand side of the inhomogeneous wave equation (1) which can be generally described by the expression

$$-\rho \frac{\partial^2}{\partial t^2}\left(\frac{\varkappa}{\rho}\frac{q}{C_p}\right),$$

and on the solution of the corresponding boundary-value problem; here, $q(t)$ is the density of the energy evolved in the medium.

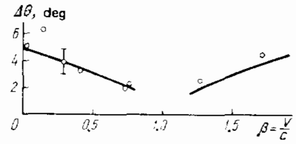

FIG. 15. Influence of the velocity of a laser beam on the half-width of the angular distribution of the radiation emitted by an opotacoustic source in the vertical plane along the direction of motion.[48]

If we assume approximately that

$$\varkappa = \varkappa_0(T_0) + \varkappa_1(T_0)(T - T_0) + \ldots,$$

and postulate that ρC_p is a constant, the solution of the boundary-value problem for the modified wave equation can be represented in the form[50]

$$p(t) = p_1(t)\left(1 + \frac{\varkappa_1}{\varkappa_0}\frac{q}{\rho C_p}\right), \tag{11}$$

where $p_1(t)$ is the solution of the corresponding boundary-value problem for Eq. (1). The quantity

$$N = \frac{\varkappa_1}{\varkappa_0}\frac{q}{\rho C_p}$$

can be regarded as a measure of the influence of the temperature-induced nonlinear effects. For example, in the case of water this quantity becomes very large in the vicinity of the point $T_0 = 4\,^\circ$C, when $\varkappa_0 = 0$.

The above expression was used to calculate numerically the acoustic signals excited in water by neodymium laser pulses of duration $\tau = 4\times 10^{-5}$ sec with q = const. The calculation was made for three values of the initial temperature of the medium $T_0 = 38$, 11, and 4 °C.

The results of a calculation of the shape of the acoustic signal at the point of observation located on the z axis (in the direction of the laser beam) in the far-field zone are plotted in Fig. 17 for different values of the nonlinear parameter N. We can see that for low values of N the nonlinear corrections are small Fig. 17(a) and as the parameter N increases, the role of

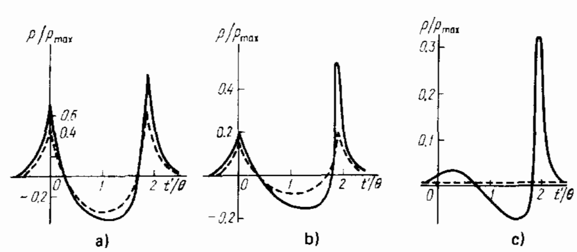

FIG. 17. Form of air acoustic signal in the far-field zone for different values of the nonlinear parameter[50]: a) $N = 0.08$; b) $N = .35$; c) $N = \infty$ ($T_0 = 4\,^\circ$C). The ordinate gives the pressure in relative units and the abscissa represents the dimensionless time. The dashed curves represent the results of the linear theory.

the nonlinear effects rises Fig. 17(b) and becomes dominant at high values of this parameter Fig. 17(c) .

The influence of nonlinearity can be described qualitatively as an increase in the "tail" part of the acoustic signal. This is due to the fact that as the energy is evolved, the volume thermal expansion coefficient of the medium increases and, therefore, the generation of sound becomes more efficient at the end of a laser pulse.

When acoustic signals are excited in water by "longer" laser pulses with harmonic intrapulse modulation of the intensity of light, the influence of the thermodynamic nonlinearity results in broadening of the acoustic signal and produces a second harmonic of the modulation frequency in its spectrum. It is interesting to note that because of the influence of the nonlinear thermodynamic effects it becomes possible to generate sound optically in water at temperatures close to 4 °C, when in accordance with the linear theory this effect should generally be absent. In this case the excited acoustic signal is governed by the second term of Eq. (11). The coefficient representing conversion of the optical energy into sound is given by

$$\eta_1=\frac{1}{4}\left(\frac{\varkappa_1 c}{C_p^2}\right)^2\left(\frac{q}{\rho}\right)^3, \tag{12}$$

i.e., the conversion coefficient rises proportionally to the cube of the energy density.

However, the value of this coefficient is usually small. For example, in the case of water at $T_0 = 20$ °C and for $q = 40$ J/cm^3 we have $\eta_1 = 10^{-4}$. For comparison, it should be noted that the efficiency of the linear mechanism of thermal generation of sound described by the first term of Eq. (11) is

$$\eta\approx\frac{1}{4}\left(\frac{\varkappa c}{C_p}\right)^2\frac{q}{\rho}. \tag{13}$$

B. Hydrodynamic nonlinearity. Parametric optoacoustic sources of sound

The formation of an acoustic signal is influenced not only by the nonlinear thermodynamic effects described above but also by the hydrodynamic nonlinearity, i.e., the nonlinearity of the hydrodynamic equations. It is manifested by the fact that an increase in the energy density of optical radiation evolved in the medium increases considerably the rate of expansion of the heated region of the medium and the intensity of the excited sound. When the sound intensity is high, nonlinear acoustic phenomena become active and they gradually *deform the profile of the excited wave as it propagates* because the different parts of the wave profile move at different velocities. This change in the profile alters the steepness of the leading edge of a compression pulse and may give rise to a weak shock wave. If p' is the peak value of the pressure in an acoustic compression pulse and l is the length of the pulse, then significant nonlinear distortions in the pulse (under planar problem conditions) occur at distance of the order of $L = l\rho c^2/\varepsilon p'$, where ε is the nonlinear parameter of the medium.[52] Under conditions typical of the thermal mechanism of optical generation of sound in liquids this distance, representing the scale of the development of the hydrodynamic nonlinearity, is usually large compared with the dimensions of the region of generation of sound (i.e., of the region of heat evolution), i.e., in other words, we always have $L \gg \mu^{-1}$. Therefore, nonlinear acoustic effects have practically no influence on the process of the optical excitation of sound in this case. On the other hand, in the region where the changes in the wave profile caused by the acoustic nonlinearity become significant the specific characteristics of the optical excitation of sound are no longer important. The problem thus splits in a natural manner into two stages: (a) calculation of the optical generation of sound using the theory of the thermal mechanism; (b) calculation of the nonlinear evolution of the newly created acoustic signal. The process of nonlinear propagation of high-power acoustic signals has been investigated thoroughly.[53] For example, it is known that in a planar one-dimensional problem a compression pulse of arbitrary shape becomes triangular after traveling a distance of the order of L and a weak discontinuity appears in its leading edge.[52] The width δ of the discontinuity is proportional to the dissipative coefficient $\nu = (2c^3/\omega^2)\alpha_s$, where α_s is the coefficient of the absorption of sound of frequency ω, and it is inversely proportional to the peak pressure p' in a compression pulse

$$\delta=\frac{\rho c\nu}{\varepsilon p'}. \tag{14}$$

Clearly, the nonlinear acoustic effects during propagation of a compression pulse are important if the width of the discontinuity is small with the spatial length of the pulse l, which yields the following condition for the manifestation of the nonlinear acoustic effects:

$$\frac{\varepsilon p' l}{\rho c\nu}\gg 1. \tag{15}$$

In addition to the multistage approach described here, one can give a theoretical description of the nonlinear acoustic effects in the optical generation of sound also in one special case when a plane wave is excited in one direction. This case may be realized, for example, when sound is generated by a scanning laser beam and the region of sound generation moves at a sonic velocity. The process of generation of a finite-amplitude plane wave is then well described (in the quadratic approximation) by the inhomogeneous Burgers equation.

The inhomogeneous Burgers equation (in the one-dimensional case), the multistage method, and the Zabolotskaya–Khokhlov equation (in the two-dimensional case) can be used to make calculations concerning the so-called parametric optoacoustic sources of sound.

In fact, if two laser beams of modulated intensity act simultaneously on the boundary of a light-absorbing medium and the modulation frequencies ω_1 and ω_2 are different, so that $\omega_1 > \omega_2$, and if the condition (15) is satisfied and strong high-frequency acoustic vibrations are generated in the light-absorption region and at distances of the order of L from it, then at considerable distances $r \gg L$ in the far-field zone we may observe acoustic vibrations of the difference frequency $\Omega = \omega_1 - \omega_2$, exactly as in the case of interaction of two high-intensity ultrasonic beams excited by a traditional method. The acoustic vibrations at the sum frequency $\omega_1 + \omega_2$ are like the vibrations of frequencies ω_1 and ω_2, damped out because of the absorption of sound in the liquid.

This represents optoacoustic parametric generation of low-frequency sound as a result of optical excitation of high-frequency ultrasonic vibrations.

A parametric optoacoustic source of sound may be constructed when an optically absorbing medium is acted upon by optical radiation whose intensity is modulated simultaneously at high and low frequencies. Then, the high-frequency ultrasonic vibrations generated in such a medium (for which, in particular, one can satisfy easily the optimal conditions for the optical generation of sound) acts as a pump in a parametric optoacoustic source of sound in whose far-field zone there are low-frequency acoustic waves. We shall conclude by noting that some of the topics discussed above are presented in Refs. 54–56.

10. OPTICAL GENERATION OF SOUND. STRONGLY NONLINEAR EFFECTS

The processes of laser generation of sound are strongly nonlinear if surface evaporation, explosive boiling, and optical breakdown of a medium take place because of strong perturbation of its equilibrium state accompanied by phase transitions and gas hydrodynamic phenomena. It should be stressed that because of the complexity of the physical effects which appear as a result of action of high-power laser radiation on matter, the acoustic aspects of these phenomena have not been investigated thoroughly and further detailed studies are needed. However, the investigations already carried out and published are sufficient for developing a qualitative description and for obtaining, in some cases, quantitative estimates for the processes of generation of sound when the density of the energy evolved in a medium is very high. The effective approaches are opposite to the acoustic theory given above and they are based, in particular, on the approximate theory of strong explosions.

A. Excitation of sound by surface evaporation of matter

The generation of sound due to the evaporation of matter becomes significant when the temperature of a medium approaches the boiling point as a result of absorption of optical radiation. If this occurs at the end of a laser pulse, an acoustic signal caused by evaporation appears in the form of an additional pressure peak in the "tail" part of the signal excited as a result of thermal expansion of the medium.[20]

An increase in the density of the energy evolved increases this maximum so that it begins to dominate the whole acoustic signal.[57] The results of detailed investigations (see, for example, Refs. 58 and 59) carried out by the method of shadow high-speed photography have shown that an increase in the intensity of laser radiation acting on the surface of a liquid (water) results in rapid evaporation of the surface layer from which a jet forms in the direction opposite to the laser beam. This jet vapor enters the air space at a high velocity and generates an intensive shock wave, whereas a recoil pulse acting on the surface of the liquid generates a compression wave. The reflection of the compression wave from the free surface, which occurs after the end of the laser pulse, creates a rarefaction wave. This causes cavitation in the surface region of the liquid manifested by readily observed bubbles. The compression wave travels in the liquid at the velocity of sound. The velocity of the shock wave in air is well into the supersonic range. For example, it was established in Ref. 60 that the action of a CO_2 laser pulse of 1.67 J energy and 75×10^{-9} sec duration on the surface of water when the diameter of the illuminated spot was $2a = 0.8$ cm produced a shock wave whose velocity in air was initially $6.4c$ (here, c is the velocity of sound in air), whereas the velocity of the compression wave in water was practically identical with the velocity of sound.

Already in one of the earliest investigations on the evaporation mechanism of optical generation of sound it was established that the pressure in a compression wave in water could reach several atmospheres.[61] Further investigations showed that the pressure amplitudes could be even greater.

Figure 18 shows, by way of illustration, the dependence of the maximum value (known as the peak value) of the pressure in an acoustic pulse in water on the optical radiation power in the form of CO_2 laser pulses (curve 1) measured at a distance of 3 cm from the sound generation region.[60] The duration of the laser pulse (measured at mid-amplitude) was 3.5×10^{-7} sec and the laser radiation was focused by a lens on the surface of water. The same figure (curve 2) shows the dependence of the amplitude of a signal excited as a result of the thermal mechanism when the laser beam was not focused. Clearly, in the case of the evaporation mechanism the peak pressure is proportional to $W - W_{th}$, where W is the laser pulse power and W_{th} is the threshold value of the laser pulse power corresponding to the onset of the evaporation of water.

This form of excitation of sound by evaporation caused by laser radiation is observed as long as the volume density of the optical energy deposited in a substance remains below a certain critical value at which optical breakdown takes place in the vapor formed from the substance. Such optical breakdown

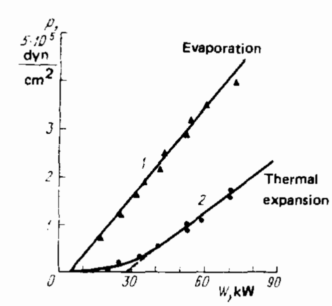

FIG. 18. Dependences of the peak pressure in a compression wave traveling in a liquid on the power of a laser pulse[57]: 1) evaporation excitation mechanism; 2) thermal expansion mechanism.

occurs, for example, when CO_2 laser radiation of $\geqslant 10^8$ W/cm^2 intensity interacts with the surface of a nonconducting liquid, and also when optical radiation of 10^6-10^7 W/cm^2 intensity is incident on the surface of a metal.

Optical breakdown in the vapor of a substance being evaporated produces a plasma which absorbs light strongly and generates a detonation wave traveling opposite to the laser beam and acquiring energy as a result of the absorption of light at the wave front. These phenomena have been discussed in detail in numerous papers on the physics of action of intense laser radiation on metals. We cannot deal with this subject in detail, and we shall only mention (by way of example) investigations of this kind reported in Refs. 62 and 63.

We can thus distinguish three regimes of laser excitation of sound when matter is evaporated and they occur in increasing order of the intensity of laser radiation: (a) weak evaporation when the density of the energy evolved in a given substance is close to the heat of evaporation; (b) rapid evaporation (explosive boiling), when the density of the evolved energy is considerably higher than the heat of evaporation, but optical breakdown of the vapor does not occur; (c) plasma formation when the intensity of laser radiation is so high that optical breakdown of the evaporation products takes place and a plasma absorbing laser radiation and screening the surface of the substance is formed.

This division into regimes of laser generation of sound during evaporation of matter is very arbitrary because the transformation of a condensed substance into a gas under the action of laser radiation is characterized by a combination of complex nonlinear phenomena, such as the formation of an overheated liquid, thermal instability of the evaporation front, surface and volume evaporation mechanisms due to the growth of unstable nuclei, possible transformation under the influence of pulsed heating of an overheated liquid to a metastable state in a surface layer, all of which result in considerable changes in the thermophysical parameters of the medium in question. Studies of these effects are important in the physics of nonresonant interaction of high-power laser radiation with matter.[64]

It is very difficult to provide a theoretical description of the evaporation mechanism of the optical generation of sound. Nevertheless, some estimates can be obtained on the basis of simple models.

Approximate calculations of the generation of sound are relatively simple at very high laser radiation intensities when optical breakdown takes place in the products of evaporation and a plasma jet is formed. In this case the phenomena observed are due to the absorption of laser radiation in a plasma and formation of a detonation wave traveling opposite to the laser beam, when a region of high pressure acting on the surface of the illuminated liquid forms behind the front of the wave.

For example, the pressure on the surface and the mechanical momentum transferred to a metal target (liquid) as a result of action of a high-power laser pulse under conditions of plasma evaporation were calculated in Ref. 62.

The initial stage of the process when the problem could be regarded as planar and one-dimensional was considered. Well-known relationships were used to find the velocity of a detonation wave and the pressure behind its front. The flow of the gas between the front of the wave and the surface of a liquid (or target) was taken to be isentropic and, assuming that the pressure behind the front of the wave was known, the solution of the one-dimensional gasdynamic problem was used to find the pressure acting on the surface of a liquid. A change in this pressure occurring after the end of a laser pulse was calculated from the well-known self-preserving solution of the initial problem of planar explosion.

The following relationships were obtained for the pressure p_{s1} on the surface of a liquid or a target during the action of laser radiation of duration τ and intensity J:

$$p_{s1}=\frac{\rho_0}{\gamma+1}\left(\frac{\gamma+1}{2\gamma}\right)^{2\gamma/(\gamma-1)}\left[2(\gamma^2-1)\frac{J}{\rho_0}\right]^{2/3} \qquad (16)$$

and after the end of such a laser pulse:

$$p_{s2}=p_{s1}\left(\frac{\tau}{t}\right)^{2/3}, \qquad (17)$$

where ρ_0 is the equilibrium density and γ is the ratio of the specific heats of the gas in front of the detonation wave.

The mechanical momentum received by the target is

$$M=6\left(\frac{\gamma+1}{2\gamma}\right)^{2\gamma/(\gamma-1)}\frac{\gamma^2-1}{(\gamma+1)^{3/2}}\left(\frac{\rho_0}{p_0}\right)^{1/2}J\tau\sim\frac{J\tau}{c}, \qquad (18)$$

where p_0 is the equilibrium pressure and c is the velocity of sound in a gas in front of the detonation wave.

Calculations carried out on the basis of Eq. (18) are in agreement with the experimental data. For example, in the case of laser radiation of 10^8 W/cm^2 intensity incident on metal targets it was found experimentally that the mechanical momentum received by the target (per unit energy of the laser pulses) was 1.1 dyn·sec $\times$J^{-1}. The corresponding theoretical value was 1.08 dyn·sec·J^{-1}.

An approximate calculation of the pressure acting on the surface of a liquid in the field of laser radiation of high intensity ($J\approx 10^8$ W/cm^2) was made in Ref. 66 for the case of strong evaporation of the liquid, but for conditions such that no plasma was formed. The time dependence of the pressure on the surface of the liquid was found using the theory of an explosive wave and also the results of Ref. 65. It was assumed that the laser pulse energy was absorbed completely by the shock wave. The law governing the motion of the front of the shock wave in air was first determined. Three stages were distinguished: motion of a plane shock wave of variable energy which occurred until the distance traveled by the wave exceeded the radius of the laser beam; motion of a spherical shock wave of variable energy which continued until the laser pulse ended; motion of a spherical shock wave of constant energy after the end of the laser pulse. The calculations showed that during the first stage of the process the pressure on the surface of the liquid could be found from the expression

$$p_s=\left(\frac{\gamma+1}{2\gamma}\right)^{2\gamma/(\gamma-1)}\left(\frac{\rho_0}{\gamma+1}\right)\left(\frac{Jc^2}{p_0}\right)^{2/3}\approx 0.16\left(\frac{Jc^2}{p_0}\right)^{2/3}, \tag{19}$$

where—as above—ρ_0, p_0, and c represent, respectively, the equilibrium density, pressure, and the velocity of sound in the atmosphere; γ is the ratio of the specific heats behind the front of the shock wave (in the calculations it was assumed that $\gamma = 1.2$).

The corresponding expression for the second stage was

$$p_s=0.16\rho_0\left(\frac{Jc^2\tau}{p_0}\right)^{2/3}(t)^{-2/3}, \tag{20}$$

whereas during the last stage of the process it was found that

$$p_s=0.16\rho_0\left(\frac{Jc^2\tau}{p_0}\right)^{2/3}t_3^{8/5}t^{-6/5}, \tag{21}$$

where

$$t_3=\left(\frac{a^3p_0}{Jc^2\tau}\right)^{1/2}. \tag{22}$$

The results of a calculation of the pressure on the surface of water acted upon by a CO_2 laser pulse are plotted in Fig. 19. In these calculations it was assumed that $J_1 = 10^7$ W/cm², $J_2 = 10^8$ W/cm², $\tau = 10^{-6}$ sec, $2a$ = 30.48 cm, p_0 = 0.1 MPa, and $c = 3.4 \times 10^4$ cm/sec.

The maximum pressure on the surface of water was calculated to be 4.8 and 22.2 MPa for the two values of the density of laser radiation adopted in the calculations.

More rigorous calculations allowing for the formation of a surface Knudsen layer in the process of evaporation can be found in Ref. 67. The results of numerical calculations[67] showed, in particular, that incidence of laser radiation of 10^6 W/cm² intensity on an aluminum target which was in air at atmospheric pressure should produce a pressure of 2 MPa on the target.

At moderate laser radiation intensities, more exactly, at such densities of the optical energy deposited in a liquid which do not cause explosive boiling of a liquid, i.e., in the case of weak evaporation, we may ignore the gasdynamic effects near the surface of the liquid subjected to laser radiation and reduce the problem to a calculation of the thermal phenomena of heating and evaporation of a condensed medium. However, one then encounters a very complex phenomenon of the liquid–vapor phase transition under the action of laser radiation. Since the depth of penetration of laser radiation into a liquid is usually greater than the thickness of the layer from which vapor is formed, the process of evaporation is governed by the transfer of energy by heat conduction from the layer where the bulk evolution of energy takes place. The nature of this process may be complicated by a number of factors. They include, in particular, the possibility of conversion by pulsed heating of an overheated liquid in the surface layer to a metastable state and this may alter greatly the thermophysical parameters of the medium. If during the existence of a metastable state one can expect spinodal singularities of thermophysical parameters, the attainment of steady-state evaporation becomes a nonmonotonic process. Volume absorption in the surface layer of a liquid accumulates energy to a level exceeding that necessary to maintain steady-state surface evaporation. This excess energy is lost by a brief increase in the temperature of the surface manifested by a peak in the evaporation pressure curve (Refs. 68 and 69).[2)]

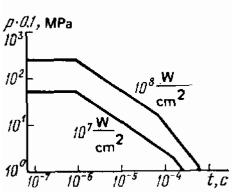

FIG. 19. Time dependences of the pressure on the surface of water for different intensities of CO_2 laser radiation in the form of pulses of 10^{-6} sec duration.[66]

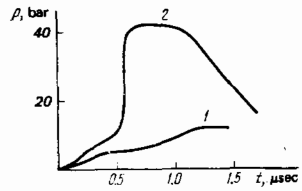

FIG. 20. Time dependences of the pressure on the surface of a liquid[69]: 1) at a low surface density of the evolved energy 3.0 J/cm² (weak evaporation); 2) at higher energy density of 3.5 J/cm².

This is illustrated in Fig. 20 which gives the results of a numerical calculation[69] of the process of evaporation under the action of laser radiation in which an allowance was made for changes in the thermodynamic parameters. At low densities of the evolved energy E = 3.0 J/cm² the evaporation pressure is described by a smooth curve 1. An increase in the density of laser radiation alters the pattern. For E = 3.5 J/cm² the curve 2 representing the time dependence of the evaporation pressure exhibits a sharp peak.[3)] The magnitude of this peak depends strongly on the intensity of laser radiation: a twofold increase in the intensity of light results in an almost twentyfold increase in the peak pressure. These numerical results are supported qualitatively by the experimental data.[70] Moreover, the presence of a temperature maximum in the surface layer of a liquid, where volume absorption of light takes place, may result in the development of a thermal instability of the evaporation front. Finally, the nature of evaporation of the liquid may change considerably as a result of the development of volume boiling effects.

As pointed out earlier, a very important characteristic of the laser generation of sound in liquids is the

[2)] However, generally speaking, because of the thermal instability of the evaporation front it is unlikely that an overheated liquid in a surface layer goes over to a metastable state.[64]

[3)] The distance traveled by light μ^{-1} is in this case very small compared with the characteristic dimensions of the problem and we can assume that the laser radiation energy is evolved on the surface of a liquid.

coefficient representing the conversion of the optical energy into sound. An estimate of the conversion coefficient under conditions of strong evaporation of a substance caused by laser radiation can be obtained using a somewhat modified theoretical model,[67] as has been done in Ref. 71.

The expression for the conversion coefficient is

$$\eta_2=\frac{[\gamma(\gamma+1)]^{2/3}(\gamma_n-1)^{4/3}p_0^{2/3}J^{1/3}}{\rho_1 c_1 c^{4/3}}, \tag{22'}$$

where p_0 and c are the equilibrium pressure and the velocity of sound in air: γ is the ratio of the specific heats in air; γ_n is the ratio of the specific heats of the vapor; ρ_1 and c_1 are the density and velocity of sound in the liquid.

We shall now give numerical estimates as an example. Let us assume that a CO_2 laser pulse of $J = 10^8$ W/cm^2 intensity is incident on the surface of water which is in contact with atmospheric air at $p_0 = 0.1$ MPa. It is then found that the proportion of the energy carried away by an acoustic wave is approximately $\eta_2 = 10^{-2}$ of the laser pulse energy.

The pressure acting on the surface of water can in this case be determined from

$$p_s=\frac{[\gamma(\gamma+1)]^{1/2}(\gamma_n-1)Jp_0^{1/2}}{\gamma_n c}. \tag{23}$$

Knowing the pressure acting on the surface of a liquid over an area equal to that of the illuminated spot, we can determine the acoustic signal excited in the liquid from the relationship

$$p=\frac{a^2}{2rc_1}\frac{p_s}{\tau}. \tag{24}$$

If we assume that in the above example the radius of a laser beam is $a = 1$ cm and the duration of a laser pulse is $\tau = 10^{-5}$ sec, then at a distance of $r = 10^2$ cm we obtain $p = 10^5$ Pa.

The above estimates are in order-of-magnitude agreement with the experimental results of Ref. 59.

B. Excitation of sound in optical breakdown of a liquid

When laser radiation is focused in the interior of a liquid, the density of the evolved energy may become so high that the liquid boils explosively in the focal region forming a rapidly expanding vapor-filled cavity and generating a compression wave.[6] At still higher intensities of optical radiation and densities of the evolved energy we can expect optical breakdown accompanied by a shock wave.[73-75] The phenomena can be described as follows. When a certain threshold intensity is exceeded, microexplosions take place in the focal region and cavities filled with a luminous plasma are formed. Laser radiation is absorbed in a dense plasma and this deposits additional energy in the cavities. The high pressure expands the cavities generating a shock wave. At the end of a laser pulse and energy evolution the gas in the plasma cavities cools, the emission of radiation from the plasma ceases, and a bubble which undergoes several pulsations is formed.

An important feature of the optical breakdown is its threshold nature. The threshold intensity of light depends on the properties of a liquid. It has been established[76] that the threshold intensity is determined by the presence of solid microparticles in a liquid, particularly particles of soot always present in the atmosphere and therefore present in liquids (water) in contact with the atmosphere. Such solid particles absorb light and are heated to temperatures of the order of 10^4 °K, which corresponds to the first ionization of the atoms and formation of a dense plasma. Light is absorbed strongly in the plasma, which causes further heating and creates a plasma cavity.

FIG. 21. Photographs of a laser spark in water from which impurities were removed.[78]

When the intensity is slightly higher than the threshold and a liquid is sufficiently pure, the number of such plasma cavities is small and they are usually located on the same straight line forming a characteristic chain or filament, which practically corresponds to the region of the maximum intensity of the incident radiation. When this intensity is increased, the bubbles merge forming a laser spark. A photograph of such a spark created by focusing a giant ruby laser pulse in water ($E = 0.6$ J, $\tau = 2\times 10^{-8}$ sec) is shown in Fig. 21 (Ref. 78). At moderate laser radiation powers the spark is nearly spherical in shape Fig. 21(a), whereas at high powers it becomes elongated Fig. 21(b).

If the concentration of microparticles is relatively high (unpurified liquid), an increase in the intensity of light increases the number of plasma cavities in the liquid and these cavities form a cloud of luminous bubbles. A photograph of such a cloud shown in Fig. 22 was obtained when a ruby laser pulse ($E = 0.2$ J, $\tau = 3\times 10^{-8}$ sec) was focused in tap water.[79]

Expansion of a plasma cavity or region generates a compression or shock wave. Compression waves can be detected optically and with the aid of wide-band

FIG. 22. Photograph of a laser spark in ordinary water.

piezoelectric transducers. For example, studies carried out using a shadow method[75,79] established that in the case of optical breakdown in various liquids (water, CCl_4, benzene) the shape of a compression wave was nearly spherical and its velocity in an interval of 3–15 μsec after breakdown was close to the velocity of sound. This is not surprising because the velocity of weak shock waves in a liquid is known to differ only slightly from the velocity of sound.

However, careful investigations demonstrated that during the initial stage of propagation of a shock wave created in a liquid by optical breakdown the velocity of such a wave can exceed the velocity of sound[80] and some departure from spherical symmetry may be observed. In particular, accelerated motion of a shock wave in the direction toward the source of laser radiation may be observed.[81]

The intensity of a shock wave rapidly decreases during its propagation. For example, an investigation of the initial state of propagation of a shock wave created by a giant neodymium laser pulse in water showed that 3.50×10^{-7} sec after the optical breakdown the velocity of the wave was 0.7 cm/μsec and the corresponding peak pressure was 230 kbar at a distance of $r = 2.5 \times 10^{-2}$ cm (Ref. 82). When a ruby laser pulse of $E = 9.6$ J energy and $\tau = 20\times 10^{-9}$ sec duration was focused in water, a compression wave recorded at a distance of 1.7 cm had a peak pressure of $p = 10^2$ atm (Ref. 77). Similar results were reported in Ref. 78. A dynamic form of the shadow method was used to study the profile of a compression wave excited by the breakdown of water caused by a ruby laser pulse of 0.1 J energy and 30×10^{-9} sec duration. The peak pressure at a distance of $r = 2.5$ cm was found to be $p = 21.5$ atm.

Although much experimental work has been done on the optical breakdown and prebreakdown phenomena in liquids, a theory of these phenomena as a whole and of the excitation of shock waves is still being developed. The excitation, propagation, and evolution of compression waves are described by the simplest model of the phenomena based on experimental observations and on the ideas developed in the theory of underwater explosions and pulsed electric discharges in water.[83-84]

It is assumed that the evolution of the optical energy fills a spherical cavity (bubble) with a homogeneous plasma and causes it to expand, so that the work is done on the surrounding liquid. The optical energy E evolved in a cavity is used only to increase the internal energy W of the matter inside the cavity and to carry out the work of expansion. Such factors as the emission of some part of the plasma energy in the form of radiation, inhomogeneous heating of the surface layers of the liquid, etc. are ignored. In the range of temperatures corresponding to the first ionization of atoms all these factors are negligible.[84]

The energy balance equation is

$$W+\int\limits_{\Delta V} p\,\mathrm{d}V=E(t); \tag{25}$$

here $V = (4/3)\pi R^3$ is the volume of a cavity and R is its radius.

The internal energy can be calculated from

$$W=\frac{pV}{\gamma-1}, \tag{26}$$

if we assume that the plasma in the cavity is characterized by some effective value of γ. For example, in the case of a plasma formed from water we have $\gamma = 1.26$ at temperatures corresponding to the first ionization of atoms.[84]

The pressure p can be found by solving the hydrodynamic problem of expansion of a sphere in a compressible liquid. This solution can be obtained in the Kirkwood–Bethe approximation. The essence of this approximation is the assumption that during the expansion of a cavity the enthalpy H travels at the local velocity of sound. The pressure is related to the enthalpy by the equation

$$p=A\left(1+\frac{n-1}{c^2}H\right)^{n/(n-1)}-B, \tag{27}$$

which follows from the equation of the state of a liquid

$$p=A\left(\frac{\rho}{\rho_0}\right)^n-B \tag{28}$$

and from the definition of enthalpy

$$H=\int\frac{\mathrm{d}p}{\rho}, \tag{29}$$

whereas the equation relating the enthalpy to the rate of expansion of the cavity $u = \mathrm{d}R/dt$ is[84]

$$\frac{\mathrm{d}H}{\mathrm{d}t}\frac{1}{c}\left(1-\frac{u}{c}\right)+\frac{H}{c}\left(1+\frac{u}{c}\right)=\frac{\mathrm{d}u}{\mathrm{d}t}\left(1-\frac{u}{c}\right)+\frac{3}{2}\frac{u^2}{R}\left(1+\frac{u}{3c}\right); \tag{30}$$

here, c is the equilibrium velocity of sound in the liquid; A, B, and n are constants. For example, in the case of water we have $A = 300.1$ MPa, $B = 300$ MPa, and $n = 7$.

The system (25)–(30) can be integrated numerically if we know the law of energy evolution $E(t)$ and the amplitude of a shock wave in a liquid can then be calculated.

Rough estimates of the hydrodynamic characteristics of the effect can be obtained from approximate relationships which follow from the system (25)–(30). Let us assume that a typical radius of a cavity reached at the end of energy evolution as a result of the absorption of a laser pulse of duration τ is R_0. Then, the characteristic rate of cavity expansion is $u = R_0/\tau$ and the pressure in the cavity is

$$p_1\approx\rho u^2=\rho\frac{R_0^2}{\tau^2}. \tag{31}$$

Using this expression, we find from Eq. (25) that the radius of the cavity is

$$R_0\approx\left(\frac{3(\gamma-1)}{4\pi\rho}\tau^2E\right)^{1/5}, \tag{32}$$ [4)]

whereas it follows from Eq. (31) that p_1 is described by

$$p_1=\rho^{3/5}\left(\frac{3(\gamma-1)}{4\pi}\right)^{2/5}\tau^{-6/5}E^{2/5}. \tag{33}$$

If $\gamma = 1.26$, which corresponds to the optical breakdown

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[4)] This expression is obtained on the assumption that the work done in the expansion of a cavity is small compared with the internal energy stored in the cavity (see also Ref. 84).

of water, we find that

$$p_1 = 0.33\tau^{-6/5}E^{2/5} \text{ cgs. units} . \quad (34)$$

In the case of a strong optical breakdown formed at intensities much higher than the threshold we can assume that the energy E evolved in the cavity plasma is of the same order of magnitude as the laser pulse energy.

When the threshold intensity is exceeded only slightly so that separate plasma bubbles are formed, we can assume that each bubble absorbs the incident optical energy in proportion to the cross-sectional area of the bubble, i.e.,

$$E = E' \frac{\pi R_0^2}{S},$$

where S is the cross-sectional area of the laser beam in the focusing region and E' is the laser beam energy.

The peak pressure in a shock wave can be calculated using approximate asymptotic formulas which follow from the theory of propagation of spherical shock waves based on the Kirkwood–Bethe approximation[80]:

$$p = \frac{p_1 R_0}{2} g \exp\left(-\frac{t}{\Theta}\right), \quad (35)$$

where $\Theta = 1 + (2/g)$; $g = (M^{3/2}\sqrt{2\ln(r/R_0)})^{-1}$ is a factor allowing for the nonlinear attenuation of the shock wave; $M = u/c = R_0/c\tau$ is the Mach number.

We shall now give examples of numerical estimates.[71] Under conditions corresponding to the experiments reported in Ref. 74 when the optical breakdown occurred in water under the action of a ruby laser pulse of $E = 0.1$ J energy and $\tau = 3\times10^{-9}$ sec duration it was found that the rate of expansion of a cavity was $u = 3\times10^3$ cm/sec, the pressure in the cavity was $p_1 = 10^4$ MPa, and the peak pressure in a shock wave at a distance of 3×10^{-3} m was $p = 37$ MPa. The corresponding experimental value was 50 MPa. Under the experimental conditions in Ref. 77, where ruby laser pulses of $E = 0.6$ J energy and $\tau = 2\times10^{-3}$ sec duration were used, it was found that $u = 3.5\times10^3$ m/sec, $p_1 = 1.2\times10^4$ MPa, and $p = 13$ MPa at a distance of $r = 1.7\times10^{-2}$ m. The experimental value was $p = 10$ MPa.

It should be noted that the above formula (35) is applicable only sufficiently far from the breakdown region where the following condition is satisfied:

$$\ln\frac{r}{R_0} \geqslant \frac{1}{4M^3}. \quad (36)$$

The above expressions can be used to find the coefficient representing the conversion of the optical radiation energy into sound:

$$\eta_3 = \frac{3}{2}(\gamma - 1)\frac{1}{M\ln(r/R_0)}. \quad (37)$$

In the above examples the value of η_3 is 20–30%, which is considerably greater than the efficiencies of the thermal and evaporation mechanisms of the generation of sound.

As pointed out in Ref. 71, the model described above may be suitable for estimating the characteristics of shock waves generated as a result of explosive boiling in the interior. We should bear in mind that, in contrast to optical breakdown, in the case of boiling explosion the deposition of the optical energy in a bubble ceases after it forms because light is reflected from the bubble surface so that the nature of the process is governed by the initial energy in the bubble.

11. CONCLUSIONS

We can thus conclude that the theory of the optical generation of sound has progressed greatly in recent years and this theory can describe fully satisfactorily the generation effect when the action of the optical radiation does not alter the aggregate state of a condensed medium. The processes of the optical generation of sound in liquids have been studied under the action of pulsed and cw lasers. The main laws governing the formation of acoustic signals have been established together with the relationships between the characteristics of these signals, on the one hand, and the parameters of optical radiation as well as thermodynamic, optical, and acoustic properties of liquids, on the other. The influence of ruffled surfaces of liquids on the optical generation of sound has been investigated theoretically. The optimal conditions for the laser generation of sound have been established. The characteristics of the generation of sound by a laser beam moving on the surface of a liquid have been investigated at subsonic, sonic, and supersonic velocities of the beam for an arbitrary form of modulation of the laser radiation intensity. Investigations have been made of the influence of the optical, thermodynamic, and acoustic inhomogeneities of the medium on the processes of the optical generation of sound. It has been shown that optoacoustic sources of sound or antennas can be established in liquids and these are capable of operating in a wide range of frequencies from the audio to the hypersound range. The frequency, directionality, and intensity of the sound emitted by an optoacoustic source can be controlled remotely by varying the parameters of the optical (laser) radiation.

The role of thermodynamic and hydrodynamic nonlinear effects in the processes of the optical generation of sound has been discussed for conditions such that the density of the optical energy evolved in a liquid is considerable, but still small compared with the heat of evaporation of the liquid. One of the possible applications is the possibility of constructing parametric optoacoustic sources of sound. It follows from the investigations carried out so far that the theory of a parametric optoacoustic source can be basically developed from the well-known theory of traditional parametric acoustic radiators.

The theory of the optical generation of sound when the dominant role is played by the thermal mechanism is in convincing agreement with the experimental results, which provides a basis for a rational selection of lasers for the generation of sound in liquids.

This can be summarized by stating that the published theoretical and experimental investigations of the optical generation of sound in condensed media relating to the thermal mechanism have advanced very greatly in recent years, although there are still many interesting topics to be tackled. For example, it would be desirable to investigate characteristics of the fields of optoacoustic sources in the near-field zone, to carry out detailed experimental studies of the characteristics of acoustic fields generated by laser radiation in different situations, and to make theoretical and experimental investigations of the characteristics of laser excitation of sound in solids.

On the other hand, there have been hardly any systematic and purposeful investigations of the laser generation of sound in condensed media under conditions of surface evaporation, explosive boiling, and optical breakdown. Considerable experimental data have already been accumulated but only fragmentary numerical and analytic estimates have been made relating to the processes of generation of sound under conditions of interaction of high-power optical radiation with matter. Many questions remain. For example, theoretical estimates given in Ref. 64 show that because of the thermal instability of the evaporation front, metastable states of a heated liquid cannot appear under conditions of surface evaporation. However, it is also pointed out in Ref. 64 that experiments deliberately designed to study the thermal instability of the phase boundary during evaporation of a liquid under the action of laser radiation have not yet been made. Consequently, we cannot exclude completely the possibility of formation of metastable states in the case of surface evaporation of liquids. It is not clear what is the relative importance of volume boiling and surface evaporation under the action of laser radiation on the free surface of a liquid.[64]

However, we have seen above that the nature of the processes occurring in the surface layer of a liquid under the action of laser radiation controls largely the efficiency of conversion of optical energy into sound, as well as the nature and evolution of an acoustic signal during its excitation and propagation.

The above theoretical and numerical estimates of the intensity of sound and of the efficiency of conversion of optical energy into sound during interaction of high-power laser radiation with matter have been made using very rough models. Nevertheless, these estimates agree in some cases with the experimental results. It follows that although the models are rough, they still reflect correctly the essence of the observed effects.

A rigorous analysis of the optical generation of sound can only be made by numerical solution of nonstationary equations of hydrodynamics, optics, and evaporation kinetics. It is worth pointing out interesting investigations[86,87] in which calculations have been made of the specific recoil momentum and processes of formation of a shock wave and its propagation into a medium, as a function of the intensity of laser radiation. The results of a numerical integration of nonstationary equations of hydrodynamics obtained without simplifying assumptions can be found in Ref. 87. The theoretically determined dependence of the specific recoil momentum on the intensity of laser radiation is in good agreement with the numerous experimental data on metals.

We have ignored above the role of electrostriction in the optical generation of sound. This role has been recently discussed in an independent review.[64] Moreover, we have neglected the mechanism associated with stimulated Brillouin scattering. Equally interesting are the optoacoustic effects associated with the resonance interaction of optical radiation with matter. These topics deserve separate discussion.

The author is deeply grateful to K.A. Naugol'nykh, S.G. Kasoev, L.V. Sedov, I.B. Esipov, S.V. Egerev, T.A. Dunina, and L.M. Bolotova for their helpful cooperation.

The author is greatly indebted to L.M. Brekhovskikh and L.V. Keldysh for their interest, and to F.V. Bunkin, A.I. Bozhkov, and V.G. Mikhalevich for valuable advice made in the preparation of the manuscript of the present paper for press.

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Translated by A. Tybulewicz