Phonon pumping of superconductors

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We show experimentally and theoretically that phonons generated in a superconductor, which acts as a wideband source of quasiparticles when illuminated by an intense laser beam, can excite a second superconductor, which is acoustically coupled to it, to a new nonequilibrium state which is sensitive to the form of the phonon spectrum. We establish conditions for diffusion-related instability of this phonon-irradiated superconductor.

INTRODUCTION

The effect of ultrasound on superconductors has been the subject of a number of publications.[1–5] In particular, in Refs. 1–3 and 5 there are discussions of "stimulated superconductivity" due to ultrasound (the Éliashberg mechanism[6]), while the authors of Ref. 4 investigated the possibility of realizing an "acoustic quantum generator" (AQG) at a frequency $\omega = 2\Delta^{(1)}$. A wideband source of quasiparticles ($\omega \gg \Delta$) cannot give rise to an AQG[4,7]; however, the energy distribution of nonequilibrium phonons is still of interest in this context, in that it can be used as a phonon pump.

In this work we will show, both theoretically and experimentally, that the phonons generated in a superconductor which acts as a wide-band source of quasiparticles when subjected to an intense laser beam can excite a second superconductor to a new nonequilibrium state. This leads to the appearance of a resistance that depends strongly on the spectrum of emitted phonons and is different when the phonon pump is in the normal and in the superconducting state. Theoretical calculation indicates that a superconductor irradiated by phonons can exhibit a "diffusion-driven instability." We note here that our use of the experimental methods described below permitted us to exclude from our investigation the purely thermal effects which always accompany experiments with laser irradiation (i.e., direct heating of the superconductor by the laser) with ease.[8–11]

I. EXPERIMENT

An experimental investigation of the behavior of superconductors under laser irradiation[8,10] shows that an observable resistance transition of superconducting-to-normal-metal type occurs over a wide range of powers β. This is explained as follows[7]: under nonequilibrium conditions the superconductor changes to an intermediate state, in which the superconducting and normal states coexist.

However, the heating of superconductors by laser irradiation meets with a serious experimental problem, to which considerable attention has been paid.[8,9,11] Hence, it is noteworthy that the experiments described below clearly document the non-thermal character of the effect of optical pumping: they show that a nonequilibrium state can arise, which is accompanied by phonon emission.

1. The configuration of the experimental apparatus is shown in Fig. 1. The thin-film lead samples under study were fastened to a special copper block located inside a superconducting solenoid cooled by liquid helium, so that the magnetic field was parallel to the plane of the film. The inhomogeneity of the magnetic field along the sample strip which was being exposed to radiation was less than 3%. A semiconductor thermometer and a Hall detector were placed together with the samples on the block. The method we used to measure and monitor the sample parameters in an external magnetic field is thoroughly described in Ref. 12. A light flux was introduced into the helium cryostat by means of a flexible light pipe and was trained on the sample. The radiation source was a GaAs injection laser with an output wavelength of 0.9 μm, a pulse duration 200 ns, and a power up to 500 W.

The layout for studying the samples is shown in Fig. 2. A wider and thicker source film, which was subjected to optical radiation, was placed on one side of an insulating substrate. On the other side a receiver film was placed symmetrically relative to the source; this film served as a recorder for phonons emitted by the first film. The films were deposited successively by thermal deposition on an insulating $8\times10\times0.5$ mm substrate or sapphire or silicon, in a VUP-4 assembly at a pressure no higher than 10^{-5} Torr. The sequence of the deposition did not influence subsequent results.

The thickness and rate of deposition were monitored directly during the preparation process by a quartz thickness gauge. The average rate of material deposition amounted to ~ 100 Å/sec. Final measurement of the film thicknesses was done with an MII-4 interferometer.

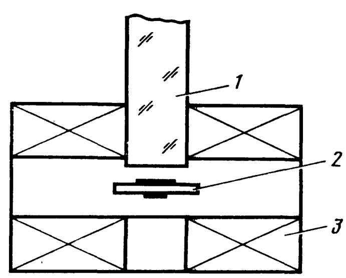

FIG. 1. Experimental Setup: 1—optical fiber, 2—sample under study, 3—solenoid.

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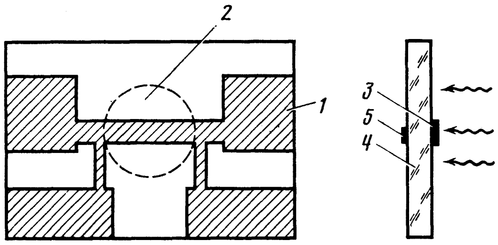

FIG. 2. Sample structure: 1—layout of sample films, 2—illuminated region, 3—source, 4—substrate, 5—receiver.

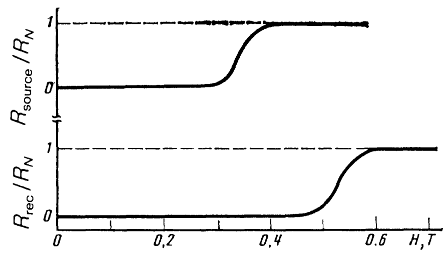

FIG. 3. Dependence of source (R_s) and receiver (R_r) resistances on magnetic field in the absence of optical pumping; $I =$ const.

We present here the results we obtained from 14 pairs of films with the following dimensions of the strips under study: source, 6×1 mm, with thickness 1000–10000 Å; receiver, 6×0.3 mm, thickness 500–1000 Å. The superconducting-film parameters were measured using the four-contact resistance method.

2. In Fig. 3 we show the characteristic dependence of the film resistance on the magnetic field. Due to the difference in film thickness, the critical magnetic field H_c^{source} for the source film was smaller than the field H_c^{rec} for the receiver. The critical temperatures were the same for both films, and the width of the transition did not exceed 0.1 K.

While the laser was trained on the sample, a pulsed current was applied synchronously with the laser pulses. Beginning with a certain value of magnetic field $H < H_c^{source}$, we consistently observed a resistive intermediate state. The current pulse appeared at the source at the instant of illumination, as confirmed by calibration using a high-speed photoreceiver. The first time at which a pulse was observed at the coincided with that of a pulse arriving at the source; its presence was connected with the direct influence of light scattered from the roughnesses sapphire substrate. When an optically opaque silicon substrate was used, this light pulse was absent. The presence of a second receiver pulse is explained by the action of phonons generated in the source by the pulse of optical radiation. The delay time of this pulse relative to the first corresponds to the ballistic propagation time of phonons in the substrate.

In Fig. 5, we show the characteristic dependence of the signal amplitude at the source on the magnetic field strength and on the laser pulse power for a number of samples. It is clear that for samples of thickness $d = 2000$–4000 Å, the critical power decreases as d increases. This implies that up to $d < 4000$ Å the inverse absorption coefficient for electromagnetic waves exceeds the thickness of the film (see Ref. 7).

Figure 6a shows a typical dependence of the amplitude of the first (optical) and second (phonon) pulses the receiver on the size of the source current pulse. The nonlinear character of this dependence suggests that the current passed through the film is close to critical. An analogous dependence on laser pulse power is shown in Fig. 6b.

A comparison of these results with Fig. 5b shows that it is impossible to estimate the critical power at which the transition to the resistive state occurs in the receiver, since the behavior of the curves is fundamentally different. This is apparently a consequence of the fact that the superconducting receiver is in a strongly inhomogeneous current state, which allows a longitudinal electric field to penetrate into the sample a macroscopic distance[13] and produce a continuous transition to the resistive state.

3. In the course of the experiments we established that the amplitude of the phonon pulse in the receiver depends on which state—normal or superconducting—the source is in: if for $H < H_c^{source}$ we drive it into the normal state by passing an above-critical current (≈ 100 mA) through it, then the phonon pulse amplitude drops with no change in the light pulse (the dashed line in Fig. 4). For magnetic fields $H_c^{source} < H < H_c^{rec}$, that is, when the source is in its normal state, passing a current of the same size through it does not lead to any change in the receiver pulse. The maximum value of this change is approximately 10% and stays constant to the limits of precision of measurement for all samples under investigation, independent of thickness.

A plot of the relative change in phonon pulse amplitude against the magnetic field for one of the samples is shown in Fig. 7. The decrease in the phonon pulse is a result of the difference between the spectra of the phonons emitted by the source superconductor in its normal and superconducting states.

II. THEORY

As was mentioned above, the appearance of pulses in the receiver is explained by the penetration of the electric

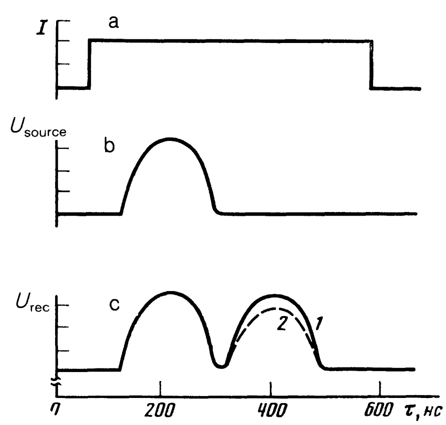

FIG. 4. Oscillograms of pulse current (a) and voltage at the source (b) and collector (c): 1—source in the superconducting state, 2—source in normal state.

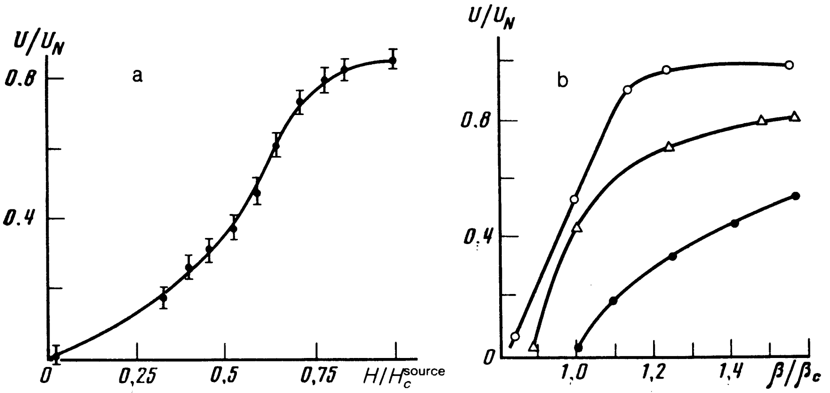

FIG. 5. Dependences of source signal amplitude: a—on magnetic field, for β, I constant; b—on optical pump power, H, I constant; ○—$d = 4000$ Å, △—$d = 3500$ Å, ●—$d = 2500$ Å; $\beta_c = \beta_c^{2500}$— critical power for the 2500 Å sample.

field to a macroscopic depth into the superconductor. The absence of a critical laser power and the peculiarities in the dependences of the pulse amplitudes on the magnetic field, current and optical pump power are all consequences of this. However, the changes in the phonon spectrum of the laser-irradiated superconductor as a function of its state, and the previously unobserved response of the other superconductor acoustically coupled to it, require some theoretical explanation.

1. Let us investigate two superconducting films acoustically coupled by a substrate, one of which (the source) is subjected to laser radiation. The distribution function $N(\omega)$ for phonons generated in the source satisfies the kinetic equation

$$N(\omega) = N_T(\omega) + \gamma J_{col}[n^1(\varepsilon)], \qquad (1)$$

where

$$N_T(\omega) = (e^{\omega/T} - 1)^{-1}, \qquad \gamma = 2\pi\lambda d\omega_D\Delta_{10}/SE_F;$$

γ is the relative probability of reabsorption and escape of phonons from the film; d, λ, ω_D, E_F and S are, respectively, the thickness, electron-phonon interaction constant, Debye frequency, Fermi energy, and velocity of sound for the source-film; $n^1(\varepsilon)$ is the quasiparticle distribution function, which satisfies its own kinetic equation[7,14]; Δ_{10} is the equilibrium value of the modulus of the source order parameter Δ_1 for $T = 0$, and J_{col} is the electron–phonon collision integral:

$$J_{col} = J_{col}[n^1(\varepsilon), N(\omega)] = \int_{\Delta_1}^{\omega-\Delta_1} \frac{d\varepsilon}{\Delta_{10}} \rho(\varepsilon)\rho(\omega-\varepsilon)\{n^1(\varepsilon)n^1(\omega-\varepsilon) - N(\omega) \times [1 - n^1(\varepsilon) - n^1(\omega-\varepsilon)]\} \times \left[1 + \frac{\Delta_1^2}{\varepsilon(\omega-\varepsilon)}\right] + 2\int_{\Delta_1}^{\infty} \frac{d\varepsilon}{\Delta_{10}} \rho(\varepsilon)\rho(\omega+\varepsilon)\{n^1(\omega+\varepsilon)[1 - n^1(\varepsilon)] - N(\omega)[n^1(\varepsilon) - n^1(\omega+\varepsilon)]\}\left[1 - \frac{\Delta_1^2}{\varepsilon(\omega+\varepsilon)}\right],$$

$$\rho(\varepsilon) = \varepsilon(\varepsilon^2 - \Delta_1^2)^{-1/2}. \qquad (1a)$$

The first term in (1a) takes into account phonon creation through quasiparticle recombination, the second term creation through scattering.

For practical film thicknesses (1000–10000 Å), $\gamma \ll 1$. This implies that nonequilibrium phonons which have not been significantly reabsorbed leave the film and thereby lose their influence on the quasiparticle distribution. Therefore, $n^1(\varepsilon)$ is a solution to the kinetic equation[7] for the equilibrium function $N_T(\omega)$. Thus, we decouple the self-consistent system of equations for phonons and quasiparticles, and dispose of the feedback, i.e., of the action of the superconductor

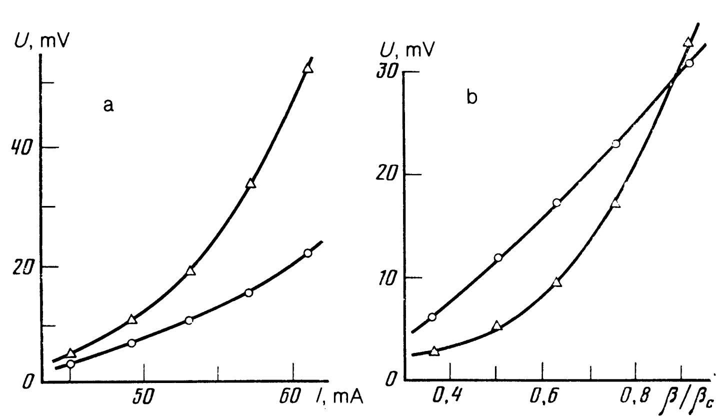

FIG. 6. Dependences of light-induced (○) and phonon-induced (△) pulse amplitudes at the collector: a—on current pulse amplitude; b—on laser optical power; β is also in units of β_c^{2500}, with H and I constant.

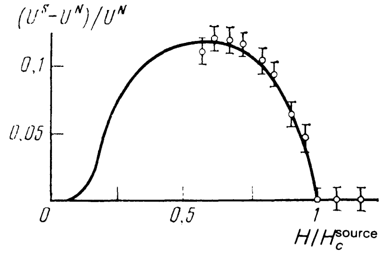

FIG. 7. Relative phonon pulse response as a function of the source state and magnetic field; U^S—phonon pulse amplitude for a superconducting source state, U^N—for a normal state. The continous curve is theoretical, the points are experimental.

phonon subsystem on the electron subsystem.

Let $T \to T_c$ for the source and receiver. We use the properties of $n^1(\varepsilon)$ in the case of optical pumping with a wide-band source of quasiparticles, i.e.,

$$n^1(\varepsilon) = (e^{\varepsilon/T}+1)^{-1} + \Delta n, \quad \Delta n \ll n^1,$$

and also the property that Δn is a function of the dimensionless pump power β. Expanding (1a) in γ and Δ_1 ($\Delta_1 \ll \Delta_{10}$ for $T \to T_c, \beta \neq 0$) and separating $N(\omega)$, we have in first approximation

$$N(\omega) = N_T(\omega) + N_1(\omega) + N_2(\omega),$$

where

$$N_1(\omega) = {}^1/_2 \gamma K(k) \theta(\omega - 2\Delta_1)$$

is the contribution from the coherent factor $\Delta_1{}^2/\varepsilon(\omega - \varepsilon)$ in the recombination integral;

$$N_2(\omega) = -\gamma F(\varphi, k)$$

is the contribution from scattering. Here,

$$\gamma = 2\gamma \frac{\Delta_1}{\Delta_{10}} f(\beta), \quad f(\beta) = \frac{1}{4} + T \frac{\partial n^1(\omega)}{\partial \omega}\Bigg|_{\omega-\Delta=0} = \beta\zeta, \quad \zeta \sim 1,$$

$$k^2 = \begin{cases} 1 - 4\Delta_1{}^2/\omega^2, & \omega > 2\Delta_1, \\ 1 - \omega^2/4\Delta_1{}^2, & \omega < 2\Delta_1, \end{cases}$$

$$\sin^2 \varphi = \begin{cases} \omega/(\omega + 2\Delta_1), & \omega > 2\Delta_1, \\ 2\Delta_1/(\omega + 2\Delta_1), & \omega < 2\Delta_1, \end{cases}$$

and $F(\varphi, k)$, $K(k)$ are elliptic integrals. Using the properties of elliptic integrals, it is easy to show that scattering processes are fully cancelled by the recombination term $\omega > 2\Delta_1$. This leads to the expression

$$N(\omega) = N_T(\omega) + N_R(\omega), \tag{2}$$

where

$$N_R(\omega) = -\gamma F(\varphi, k) \theta(2\Delta_1 - \omega). \tag{3}$$

Thus, for $\omega < 2\Delta_1$, there appears a "phonon deficit" (see Ref. 3), since the source film absorbs in this energy interval phonons that participate in the reabsorption processes (Fig. 8). This happens for the following reason: in our approximation the rates of phonon generation due to quasiparticle recombination and phonon annihilation due to scattering are equal, and are proportional to γ and Δ_1. However, phonons are generated by scattering in the whole energy range $\Delta_1 < \varepsilon < \infty$, while phonon generation due to recombination begins only when $\varepsilon > 2\Delta_1$. From this energy on, both processes cancel each other, so that there remains only the equilibrium function $N_T(\omega)$. As a result, for $\omega < 2\Delta_1$ phonon annihilation due to quasiparticle scattering leads to a relative "phonon deficit." We remark that the "phonon deficit" was predicted earlier for electromagnetic pumping at frequencies $\Omega < 2\Delta_1$ (Ref. 3).

2. Nonequilibrium phonons which leave the source excite the receiver and introduce in it the distribution function $n(\omega)$. Let us write down an equation for the quasiparticle distribution function $n(\varepsilon)$ for the receiver in the form[7,14,15]

$$I_{ph}[n(\varepsilon), N_T(\omega)] = Q_1(\varepsilon) + Q_2(\varepsilon) + Q_3(\varepsilon),$$

$$Q_1(\varepsilon) = \int_0^\infty d\xi' \rho(\varepsilon, \varepsilon')[1 - n(\varepsilon) - n(\varepsilon')] N_R(\varepsilon + \varepsilon'),$$

$$Q_2(\varepsilon) = \int_0^\xi d\xi' \rho(\varepsilon, -\varepsilon')[n(\varepsilon') - n(\varepsilon)] N_R(\varepsilon - \varepsilon'), \tag{4}$$

$$Q_3(\varepsilon) = \int_\xi^\infty d\xi' \rho(-\varepsilon, \varepsilon')[n(\varepsilon') - n(\varepsilon)] N_R(\varepsilon' - \varepsilon),$$

$$\rho(\varepsilon, \varepsilon') = (\varepsilon + \varepsilon')^2 (1 + \Delta^2/\varepsilon\varepsilon'), \quad \xi = (\varepsilon^2 - \Delta^2)^{1/2}.$$

Here, I_{ph} is the electron-phonon collision integral.

Near T_c,

$$n(\varepsilon) = n_T(\varepsilon) + \tilde{n}(\varepsilon), \quad \tilde{n} \ll n_T, \quad n_T(\varepsilon) = (e^{\varepsilon/T}+1)^{-1}.$$

We supplement equation (4) by the self-consistency condition[2]

$$\Delta[\alpha + \beta\Delta^2/T_c{}^2 + \chi] = 0, \tag{5}$$

$$\alpha = \frac{T - T_c}{T_c}, \quad \beta = \frac{7\zeta(3)}{8\pi^2}, \quad \chi = 2\int_0^\infty \frac{d\xi}{\varepsilon} \tilde{n}(\varepsilon).$$

Let us investigate some limiting cases.

3. $\Delta_1 \ll \Delta$. In the source term there remain the terms $Q_{2,3}(\varepsilon)$. Taking into account that the characteristic energy ε by which $\tilde{n}(\varepsilon)$ is changed is of order Δ. We obtain in lowest approximation in Δ/T (and neglecting the integral term as $T \to T_c$),

$$\tilde{n}(\varepsilon) = \bar{\gamma}\{(\varepsilon + 2\Delta_1 - \Delta^2/\varepsilon)[(\varepsilon + 2\Delta_1)^2 - \Delta^2]^{-1/2} - (\Delta_1 \to -\Delta_1)\},$$

$$\bar{\gamma} = \frac{\pi\gamma}{14\zeta(3)} \frac{\Delta_1{}^4}{T^4}. \tag{6,I}$$

We emphasize that (6) has the form of the solution for an AQG of frequency $\omega = 2\Delta$, apart from the sign.[1,2]

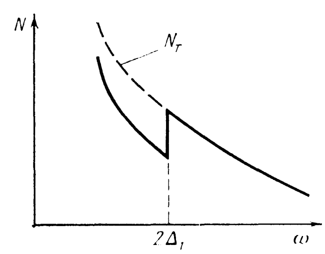

FIG. 8. Theoretical phonon spectrum for the source film.

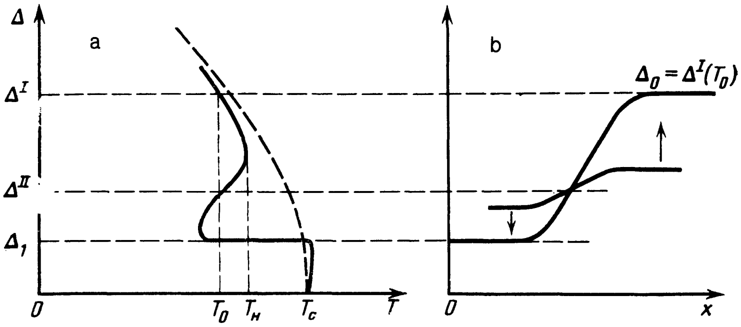

FIG. 9. Dependence of the receiver order parameter on temperature (a) and coordinate (b) (layer-solution).

Substituting (6) into the expression (5) for χ, we obtain

$$\chi=\frac{4\bar{\gamma}\Delta_1}{\Delta+\Delta_1}\left\{K(k)-\frac{\Delta^2}{\Delta_1^2}\left[\Pi(\sigma^2,k)-K(k)\right]\right\},$$

$$k=(\Delta-\Delta_1)/(\Delta+\Delta_1),\quad \sigma=\Delta_1/(\Delta+\Delta_1).$$

Expanding the elliptic integrals K and Π in terms of $\Delta_1/\Delta\ll1$, we have

$$\chi=4\bar{\gamma}\Delta_1/\Delta. \tag{7}$$

From the self-consistency equation we find that finite nontrivial solutions to (5) exist up to δ_H (T_H) values satisfying the relation

$$\delta_H^2=\frac{4}{27\beta}\left(\frac{T_c-T_H}{T_c}\right)^3,\quad T_H<T_c,$$

while the order parameter is not single-valued (Fig. 9a):

$$\Delta^{\rm I}=\frac{2}{\sqrt{3}}\Delta_0(T)\cos\frac{A}{3},\quad \Delta^{\rm II}=-\frac{2}{\sqrt{3}}\Delta_0(T)\cos\left(\frac{A}{3}+\frac{\pi}{3}\right), \tag{8}$$

$$\cos A=-\delta/\delta_H,\quad \Delta_0(T)=T_c(-\alpha/\beta)^{1/2}.$$

Here $\delta=4\bar{\gamma}\Delta_1/T_c$ is a monotonic function of Δ_1, playing the role of the dimensionless power of the source of nonequilibrium phonons.

4. $|(\Delta_1-\Delta)/\Delta|\ll1$. The narrow-source conditions,[7,14] hold for $\tilde{n}(\varepsilon)$, and recognizing that in the coherent factors the sign is different from the case of tunneling injection and optical pumping (this allows us to neglect the terms $Q_{2,3}(\varepsilon)$ as $T\to T_c$), we obtain the equation

$$\tilde{n}(z)\Gamma(T)+[N_T(z)+n_T(z)]\bar{n}$$
$$=-{}^1/_8\pi\gamma(\Delta/T)\tilde{\delta}^{1/2}(1-z)^{1/2}\theta(z)\theta(1-z),$$
$$z=\frac{\varepsilon-\Delta}{2\Delta\tilde{\delta}},\quad \tilde{\delta}=\frac{\Delta_1-\Delta}{\Delta},\quad \Gamma(T)=\frac{7\zeta(3)}{16}\frac{T^3}{\Delta^3}, \tag{9}$$
$$\bar{n}=\frac{1}{\Delta}\int_0^\infty\frac{\varepsilon d\varepsilon}{(\varepsilon^2-\Delta^2)^{1/2}}\tilde{n}(\varepsilon).$$

Neglecting the integral term as $T\to T_c$, we find from (5) for small deviations of Δ from the equilibrium $\Delta_0(T)$, in light of (9),

$$\frac{\Delta-\Delta_0(T)}{\Delta_0(T)}=\frac{8\pi^2\gamma}{49\zeta^2(3)}\frac{\Delta_0^2(T)}{T^2}\delta(T)\theta[\delta(T)], \tag{10}$$
$$\delta(T)=[\Delta_1-\Delta_0(T)]/\Delta_0(T).$$

From (10) it is clear that stimulated superconductivity occurs at $\Delta_1\gtrsim\Delta$ (Fig. 9a).

5. Let us investigate the inhomogeneous state of the receiver film. In this case, the equation for Δ takes at $\xi_0\gg L$ (ξ_0 is the coherence length, L is the diffusion length) the form ($\Delta\gg\Delta_1$)

$$\xi_0^2\frac{d^2\Delta}{dx^2}=\Delta\left[\alpha+\beta\frac{\Delta^2}{T_c^2}+\frac{\delta T_c}{\Delta}\right],\quad \left.\frac{d\Delta}{dx}\right|_{\pm l/2}=0 \tag{11}$$

(l is the sample length). The only stable inhomogeneous solution of (11) is that for layers[7]; and realized for $\delta_0(T_0)$ with

$$\delta_0^2=\frac{2}{27\beta}\left(\frac{T_c-T_0}{T_c}\right)^3,\quad \delta_0=\frac{\delta_H}{2},\quad \frac{T_c-T_H}{T_c-T_0}=2^{1/3}.$$

This solution separates regions with different values of the order parameter $\Delta=\Delta_0$ and $\Delta\sim\Delta_1\ll\Delta_0$; for $\Delta\ll\Delta_1$ it takes the form (see Fig. 9b)

$$\Delta=\Delta_0\{1-3[\mathrm{ch}\,((-\alpha)^{1/2}\xi_0^{-1}(x-x_0)+C)+2]^{-1}\}, \tag{12}$$

where $\cosh C=4$, x_0 is the coordinate of the phase boundary, and $\Delta_0=-3\delta_0T_c/\alpha_0$, $\alpha_0=(T_0-T_c)/T_c$.

For $\delta\neq\delta_0$ ($T\neq T_0$), the boundary separating the phases moves with velocity $v_c\sim(\delta-\delta_0)/\delta_0$, so for $\delta<\delta_0$ it changes the sample into the homogeneous state $\Delta=\Delta_0$ and for $\delta>\delta_0$ into a state with $\Delta\sim\Delta_1\ll\Delta_0$. For $L\gg\xi_0$ these results do not change qualitatively.

If a current flows through the receiver, then according to Ref. 13, a longitudinal electric field penetrates to a distance $l_E\sim LT^{1/2}/\Delta^{1/2}$ into the superconductor. (In the nonequilibrium situation, this quantity is significantly larger.) Thus, in the presence of a current the collector goes into a resistive state, and allowance for the stratification into phases leads to an additional resistance correction due to the regions with $\Delta\sim\Delta_1$. We emphasize that this correction obtains only for $\Delta_1\neq0$, that is, for a superconducting state of the source.

An estimate of the contribution of this effect to the resistance gives a dependence entirely in agreement with experiment (Fig. 7).

6. At $\Delta_1\ll\Delta$, the conditions for diffusion-driven instability (DDI) obtain in the superconducting receiver. According to Refs. 7 and 17, the appearance of a DDI is connected with the sign of

$$N_S=1+2\int_0^\infty\frac{\partial n}{\partial\varepsilon}d\xi.$$

If $N_S<0$ (which corresponds formally to a negative coefficient of quasiparticle diffusion), then a DDI can occur. Substituting (ε) for $\Delta_1 \ll \Delta$ in N_S, we have with an accuracy up to terms Δ^2/T^2

$$N_s = 2\int_0^\infty \frac{\partial \tilde{n}}{\partial \varepsilon} d\xi = -2\bar{\gamma}\frac{\Delta_1}{\Delta}.$$

The presence of a DDI in our case[2)] implies that thanks to the anomalous diffusion of quasiparticles, the stratification into phases with differing values of the order parameter can occur even without priming inhomogeneities (phase nucleation centers), that is, it can occur in an absolutely pure superconductor.

III. DISCUSSION OF RESULTS

The experimental results we have obtained regarding the dependence of the relative change of the phonon pulse amplitude (Fig. 7) on the magnitude of the magnetic field are in complete agreement with the theoretical calculations, if we take it into account that a change in the magnetic field changes the relative magnitudes of the order parameters of the acoustically-coupled layers because of the difference in their thicknesses.[19] We note that for $H/H_c^{\text{source}} < 0.5(\Delta_1/\Delta \gtrsim 0.5)$ the theoretical results point also to a diminution of these effects, but no experimental points could be recorded in this region. The absolute values of the receiver signals and their dependence on the magnetic field, the current and the optical pump power are also rather well described by invoking the effect of electric field penetration into the superconductor.[13] We emphasize that this effect is sizable because of the small layer thickness and the closeness of the current value to critical.

Let us formulate the basic results of this work:

1. The experimental observation of destruction of superconductivity under the action of laser radiation is explained by an excess (i.e., relative to thermal) of excitations of the electron spectrum. This is confirmed by the difference between the spectra of phonons emitted in the normal and superconducting states.

2. The relative change in the response of the superconductor to phonon pumping was calculated theoretically. The results agree completely with the experimental data.

3. It was shown theoretically that for a certain type of phonon-pumping spectrum there can arise a diffusion-driven instability in the superconductor, which contributes to the stratification of the superconductor into phases with different values of the order parameter.

4. It was calculated that when two acoustically-coupled superconductors have order parameters whose values are close to one another, superconductivity can be stimulated; the stimulation mechanism differs from that of Eliashberg, since here recombination processes play the principal role in the formation of the nonequilibrium distribution function.

5. All the results obtained here are a consequence of the "phonon deficit" which occurs for frequencies less than 2Δ.

It is necessary to note that ultrasonic pumping would lead to opposite effects[1-5]: suppression of the superconductivity at frequencies above 2Δ, and stimulation by the Eliashberg mechanism at frequencies below 2Δ.

It should also be noted that a shift in the I–V characteristic was observed in Ref. 20 for a film acted on by nonequilibrium phonons generated in a source film by an electron beam. Because an electron current gives rise in the superconductor to a nonequilibrium state analogous to that which appears under the action of optical pumping with a wideband quasiparticle source, the receiver-signal decrease observed in Ref. 20 when the source returned to the normal state can also in principle be attributed to "phonon deficit" at $\omega < 2\Delta$.

The authors express their gratitude to V. F. Elesin for constant attention to this work and for helpful discussions.

1) $\hbar = e = c = 1$.

2) A multi-valued dependence of the order parameter on the source phonon frequency and conditions for a DDI were obtained also in Ref. 2, while in Ref. 18 it was shown that a DDI can occur in superconductors with a strong electron-phonon interaction.

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