Destruction of superconductivity by optical radiation and nonequilibrium resistive states in films of the high-T_c superconductor $YBa_2Cu_3O_{7-x}$

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The nonequilibrium resistive response of thin films of the high-T_c superconductor $YBa_2Cu_3O_{7-x}$ to laser light has been studied experimentally and theoretically. Two relaxation times can be distinguished in the resistive response of the films: a "short" time $\tau_s \leqslant 10$ ns and a "long" one $\tau_l \gg 10\,\mu$s. The latter is associated with a heating of the superconductor's lattice. The relative sizes of the thermal and nonthermal components of the resistive response are studied for various values of the parameters (the temperature, the intensity and length of the laser pulses, the dimensions of the samples, and the transport current). The observed effects are interpreted in a model with a spatially nonuniform distribution of the order parameter $\Delta(\mathbf{r})$ in the films and with penetration of an electric field into regions of pronounced nonuniformity as a result of "hot" quasiparticles.

1. INTRODUCTION

Experiments on the destruction of superconductivity by optical radiation have been carried out for a long time (Refs. 1–4, for example). Testardi[1] observed a nonthermal response of a superconducting film at a temperature $T < T_c$, where T_c is the superconducting transition temperature, to external optical radiation. "Nonthermal" here means an effect which is unrelated to a heating of the lattice of the superconductor. The "response" of the system is understood here as the appearance, upon the application of optical radiation, of a finite voltage U across a superconducting film through which a direct current $j < j_c$ is flowing (j_c is the critical current). The relaxation times of this voltage have been observed to be far shorter than the time scales of the lattice cooling as a result of thermal conductivity. At low light intensities, no response is observed; the response appears after a threshold is exceeded, at $I > I_{c1}$. At $I > I_{c2}$, the resistance of the sample which arises because of the light is the same as that of the sample in its normal phase (we denote the corresponding voltage by U_N). In the intensity interval $I_{c1} < I < I_{c2}$ the voltage U varies from zero to U_N. There thus exists an intensity interval in which an incomplete destruction of the superconductivity is observed. The experiments and their results are described in detail in Refs. 1 and 2. Several theoretical papers, e.g., Refs. 3–7, have been devoted to an interpretation of these experiments. The reader interested in a more detailed bibliography of the work in this area is directed to the reviews by Elesin and Kopaev[3] and Aronov and Spivak.[4]

From the way in which the experiments were carried out, we understand that the resistive response of the superconductor stems from the appearance of some deviation from equilibrium in this case. This response was attributed in Refs. 3 and 5–7 to a deviation of quasiparticles from equilibrium. In the present paper we will also assume that only the system of quasiparticles deviates from equilibrium. This approach is justified in a study of narrow bridges on thin high T_c films, in which there is a fast "ballistic" escape of nonequilibrium phonons from the sample. On the other hand, we should point out that the particular deviation from equilibrium, which is playing the major role, may be caused by other subsystems, depending on the particular physical situation. These other subsystems might be nonequilibrium phonons or a magnetic subsystem associated with the magnetic field of the current flowing through the film.

One of the basic results of Ref. 3 is a multivalued dependence of the width of the superconducting gap on the intensity I of the optical radiation. Since there are coherence factors in the collision integral representing collisions of quasiparticles with phonons, the average recombination rate of the quasiparticles falls off with decreasing value of the superconducting gap Δ at small values of Δ. This effect leads to a coherence instability,[3,5–7] which causes the superconducting gap $\Delta(I)$ to abruptly vanish when the intensity of the optical radiation exceeds a certain critical value I_{c2}. In other words, the superconducting film is in its normal phase at $I > I_{c2}$. In the light intensity interval $I_{c1} < I < I_{c2}$, on the other hand, where $\Delta(I)$ is multivalued, the normal and superconducting phases may coexist. In this "switching-wave" regime, a wave which switches the sample from its superconducting state to its normal state or vice versa propagates through the film. This is the qualitative explanation offered in Refs. 3 and 5–7 for the experiments of Refs. 1 and 2 on the optical destruction of superconductivity in thin superconducting films.

However, the mechanism discussed in Refs. 3 and 5–7 for the optical destruction of superconductivity in thin films is apparently incapable of giving a complete description of the appearance of resistive states when optical radiation is applied to high T_c superconducting films. Here are the arguments for that assertion.

The critical intensity (for the occurrence of a coherence instability) can be estimated in a straightforward way from Refs. 3 and 5–7. Complete destruction of the superconductivity should correspond to that density of quasiparticles at which Δ vanishes. From the condition for a steady state we have $R = G$, where R is the recombination rate of the quasiparticles, and G is the average rate of their photoproduction (the rate at which the quasiparticles are produced by the optical radiation). Since R is an average collision integral representing collisions of quasiparticles with phonons, we should have $R \sim \gamma N^2$, where N is the density of quasiparti-

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cles, and $\gamma \approx$const. However, the quasiparticle density which would be required for complete destruction of the superconductivity is actually determined by a phase volume at the Fermi surface with a thickness (along the energy scale) on the order of Δ_0, where Δ_0 is the superconducting gap at absolute zero.[1] We find thus $N \sim \Delta_0$, i.e., $R \sim \gamma \Delta_0^2$. The rate G of quasiparticle photoproduction, on the other hand, is proportional to the light intensity I. The condition for a steady state then leads us to

$$\gamma \Delta_0^2 \sim R = G \sim I_{c2},$$

so the critical light intensity for the destruction of superconductivity throughout the sample is proportional to Δ_0^2. The dependence $I_{c2} \propto \Delta_0^2$ shows which light intensities are required for this total destruction of the superconductivity in a high T_c superconductor, since the parameters of the superconductor (which figure in γ) other than Δ_0 are not greatly different from their values in the low-temperature superconductors. We immediately find an estimate of the critical light intensity at $T \ll T_c$:

$$\frac{I_{c2}^{\rm HTSC}}{I_{c2}^{\rm LTSC}} \sim \left[\frac{\Delta_0^{\rm HTSC}}{\Delta_0^{\rm LTSC}} \right]^2 \sim \left[\frac{T_c^{\rm HTSC}}{T_c^{\rm LTSC}} \right]^2 \sim 10^2 - 10^3,$$

where $\Delta_0^{\rm HTSC}$ is the "bulk" value (the value within the crystallites) of the order parameter in the high T_c superconductors. In the experiments of Refs. 1 and 2, however, the typical intensities at which superconductivity was destroyed were on the order of 10^3 W/cm^2. According to our estimates, the critical intensities should have been on the order of 10^5–10^6 W/cm^2 in the high T_c materials if the same mechanism for the destruction of superconductivity had been operating. In the high T_c materials, however, resistive states are known[8,9] to arise even at substantially lower light intensities (three or four orders of magnitude lower). The estimate above is valid for a uniform distribution of the order parameter and for $\Delta \sim \Delta_0$. The high T_c films which can presently be grown, however, are inhomogeneous (Ref. 10, for example), and the order parameter Δ may have different values in different parts of the film. Consequently, in interpreting experimental results one must bear in mind that the parameter Δ may be small in certain regions (and locally quasiuniform) and that the mechanism proposed in Ref. 3 may operate in these regions.

A spatial inhomogeneity appears to be an important property of the high T_c materials. It probably stems not from a "low quality" of specific samples but from the extremely small value of the coherent length ξ^* in these materials. In $YBa_2Cu_3O_{7-x}$, this length is[11] $\xi_c^* \approx 3$–4 Å along the c axis, while that in the ab plane is $\xi_{ab}^* \sim 20$–30 Å, i.e., on the order of a few interatomic distances. Whereas in low-temperature superconductors, where ξ^* is larger by several orders of magnitude, the "proximity effect"[12] prevents irregularities with a size of a few atomic length scales from causing a substantial spatial variation of the superconducting gap $\Delta(\mathbf{r})$, in the high T_c superconductors, with $\xi^* \leqslant 10$ Å, any irregularity of the material or of the superconductor or substrate with a size greater than or on the order of a few times the atomic length scale ($\geqslant \xi^*$) will lead to a substantial spatial variation of $\Delta(\mathbf{r})$.

Moreover, thin films of the high T_c materials are not "solid," in the sense that they consist instead of blocks (in the case of single-crystal films) or grains (in the case of polycrystalline films) with typical sizes $a \sim 100$–1000 Å. As a result, there are regions (crystallites) in which Δ is comparatively large. These regions occupy a large fraction of the volume of the superconductor. There are also regions, between these crystallites, in which the values of Δ are considerably smaller and are to some extent random, with a distribution function which depends on the particular conditions under which the given sample was grown.

This fact has been discussed previously in the literature. Attempts have been made (e.g., Ref. 13) to interpret the properties of high T_c superconductors in a model of a random network of intergrain Josephson junctions. It is thus incorrect to analyze the destruction of superconductivity against the background of a uniform state; it is necessary instead to consider the optical destruction of superconductivity in the case of a spatially nonuniform distribution of the size of the superconducting gap, $\Delta_0(\mathbf{r})$. A nonuniform $\Delta_0(\mathbf{r})$ profile in the high T_c superconductors might also have some new experimental consequences.

In this paper we are reporting an experimental and theoretical study of the properties of films of this sort—spatially nonuniform thin films of high T_c superconductors—when exposed to optical radiation. In Sec. 2 we describe our experimental results. In Sec. 3 we take up a theoretical model which (we believe) makes it possible to interpret these experiments.

2. EXPERIMENTAL PROCEDURE AND RESULTS

In this section of the paper we examine the results of experiments on the magnitude and temporal characteristics of the nonequilibrium resistive response of superconducting $YBa_2Cu_3O_{7-x}$ films to applied laser light at various values of the parameters, namely, the film thickness, the transport current, and the temperature.

a) Test samples and experimental procedure

In the experiments we studied the properties of $YBa_2Cu_3O_{7-x}$ superconducting films grown at our Institute by laser deposition on Zr_2O and $SrTiO_3$ substrates. The films ranged in thickness from 1500 to 6000 Å and had a block structure, with the crystallographic **c** axis perpendicular to the plane of the film. We know that the blocks (or grains) in such films have a size in the ab plane on the order of 1000 Å. These films were used to fabricate narrow bridges, ranging in width from 50 to 500 μm, with lengths of 3–8 mm. Silver contacts were deposited on the bridges for measurements by the four-probe method (Refs. 1 and 2, for example). The contact resistance at $T = 77$ K was no higher than 1 Ω. Figure 1 shows typical curves of the resistance of these bridges versus the temperature T for various values of the transport current. It can be seen from this figure that the superconducting transition temperature for the samples used in these experiments is $T_c \approx 84$–85 K, and the transition width is no greater than $\Delta T_c \approx 2$–3 K. The critical current density J_c in these bridges at $T = 77$ K is $J_c \geqslant 2 \cdot 10^5$ A/cm^2.

Figure 2 shows the layout of our experimental apparatus. A direct transport current $J \ll J_c$ was passed through the film of the high T_c superconductor with the help of a pair of current contacts. The film was housed in a cryostat. The optical radiation from a pulsed neodymium laser (with an

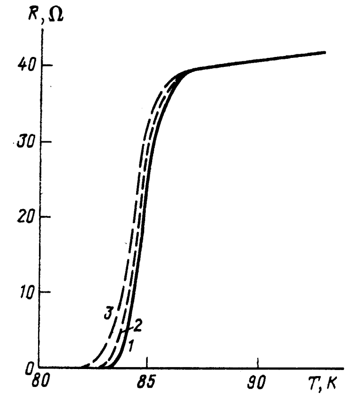

FIG. 1. Temperature dependence of the resistance of a high T_c bridge with a thickness $h = 3300$ Å at various current densities $j(\mathrm{A/cm^2})$: 1—$3.1\cdot10^2$; 2—$1.4\cdot10^3$; 3—$2.6\cdot10^3$.

output wavelength $\lambda_{\mathrm{las}} = 1.06$ μm and a pulse length $\tau_{\mathrm{las}} \approx 30$ ns) was projected onto a region $d \approx 1$ mm in diameter between the contacts. The sample was exposed to the light in such a way that the entire width of the bridge was illuminated. Since the length of the bridge was large in comparison with the diameter of the exposed region, and since the contact resistance was low, we were able to eliminate almost completely the effect of illumination of the contacts in these measurements. The potential contacts were used to measure the voltage $\Delta U = \Delta RJ$, i.e., the response of this current-carrying superconducting system to the pulse of optical radiation. This voltage was sent to an oscilloscope. For the results reported below, the power density of the laser light was a constant 1.5 kW/cm^2.

b) Experimental results

Analysis of the temporal characteristics of the voltage pulses which appeared across the potential contacts revealed two distinct characteristic relaxation times in the response of the system: a "short" time $\tau_s \leqslant 10$ ns and a "long" one $\tau_l \gg 10$ μs. The relative sizes of the thermal and nonthermal components of the response depend on the various experimental parameters and conditions, e.g., the film thickness, the temperature, the transport current, and the width of the superconducting transition, ΔT_c.

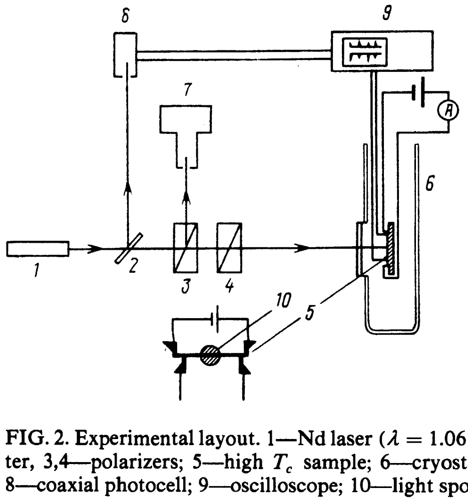

FIG. 2. Experimental layout. 1—Nd laser ($\lambda = 1.06\,\mu$m); 2—beam splitter, 3,4—polarizers; 5—high T_c sample; 6—cryostat; 7—power meter; 8—coaxial photocell; 9—oscilloscope; 10—light spot.

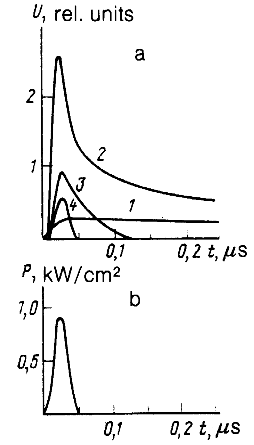

FIG. 3. a: Oscilloscope traces of the photoresponse of a high T_c film to the laser pulse shown in part b of this figure, at various temperatures. 1—100 K; 2—86 K; 3—84 K; 4—77 K. b: Oscilloscope trace of the laser pulse acting on the film.

Figure 3 shows an oscilloscope trace of the response of the system, observed as a voltage signal. The presence of an equilibrium thermal component of the response can be established experimentally in a fairly easy way (Fig. 4), from the shape of the responses to a train of several laser pulses. if the system has not completed its thermal relaxation by the time the next pulse of the train arrives, each pulse of the photoresponse is raised above the preceding pulse (Fig. 4b). A time shift of the photoresponse with respect to the beginning of the laser pulse at $T < T_c$ is evidence that a threshold (in the laser light intensity) must be exceeded for the nonequilibrium resistive state to arise. The presence of a characteristic time on the order of τ_l in the response (see the curve in Fig. 3 corresponding to $T = 100$ K) is usually linked with a bolometric (thermal) effect: the appearance of a resistive state as the result of a heating of the lattice. The time τ_l is determined by the characteristic time for heat removal from the system. The presence of the characteristic relaxation time τ_s, on the other hand, indicates the existence of a non-

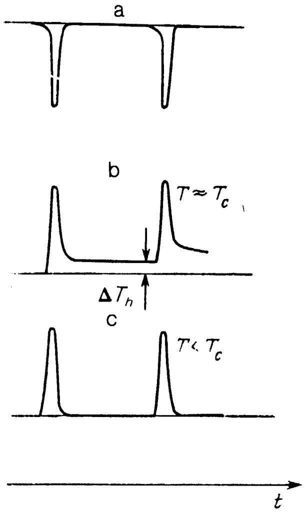

FIG. 4.

thermal mechanism for a deviation from equilibrium in this system. We will refer to this mechanism as the "nonbolometric" (or "nonthermal") mechanism for a deviation from equilibrium.

As we have already mentioned, the nonthermal response of a current-carrying superconducting film to external optical radiation was observed a fairly long time ago in low T_c superconductors (Refs. 1 and 2, for example). In the high T_c materials, however, the mechanisms for the occurrence of this deviation from equilibrium and the region of parameter values in which they operate (and are manifested experimentally) may be qualitatively different. Furthermore, the deviation from equilibrium which arises in the high T_c materials has several important distinguishing features and may lead to some interesting new results, which may be of interest for applications.

Curve 1 in Fig. 5 shows the amplitude of the response of the system to external optical excitation (i.e., the photoresponse of the system) versus the sample temperature at a transport current density $j = 4.6\cdot 10^3$ A/cm^2. Curve 2 in Fig. 5 shows for comparison the corresponding behavior of the bolometric component of the photoresponse. This curve was plotted with the help of information on the heating $\Delta(T_h)$ of the superconductor's lattice found experimentally (as in Ref. 2, at $T > T_c$) and on the experimental temperature dependence of the bridge resistance R:

$$\Delta U_{bol} \approx \Delta T_h (dR/dT) J_t.$$

It was established experimentally that at the power levels of the optical radiation used in our experiments the heating of the samples did not exceed 0.5–0.7 K.

It can be seen from Fig. 5 that the total photoresponse signal may differ substantially from the thermal signal, particularly as the temperature is lowered ($T < T_c$). It then becomes a straightforward matter to work from the total response of the system to the external radiation to find the nonthermal component of the response. Figures 6 and 7 show the amplitudes of the thermal and nonthermal components of the photoresponse versus the temperature for various values of the transport current density J_c found by this method. It turns out that the amplitude of the nonthermal

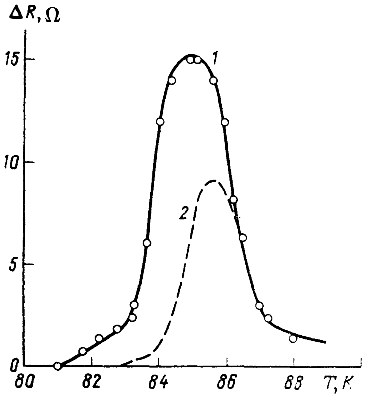

FIG. 5. The photoresponse ΔR versus the temperature of a high T_c sample with a thickness $h = 6000$ Å at a current density $j = 4.6\cdot 10^3$ A/cm^2. 1—Total signal; 2—thermal component.

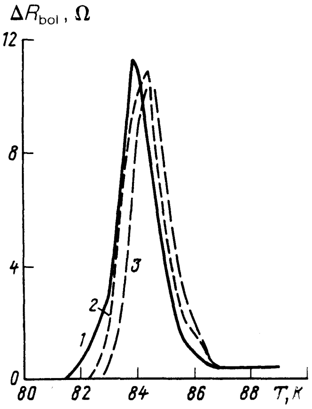

FIG. 6. Thermal component of the photoresponse ΔR versus the temperature of the sample at various current densities j (A/cm^2). 1—2.6$\cdot 10^3$; 2—1.4$\cdot 10^3$; 3—3.1$\cdot 10^2$. The bridge thickness is $h = 3300$ Å.

response is the stronger function of the transport current, increasing with increasing J_c. The nonthermal response exists at lower temperatures (Figs. 6 and 7). The highest amplitudes of both the nonthermal and thermal components are found at progressively lower temperatures as the transport current is raised.

In singling out the nonbolometric response of the system by the method outlined above, we assumed that the heating of the sample, ΔT_h, was approximately the same above and below T_c at a given laser light intensity. This assumption is legitimate if the specific heat and thermal conductivity of the substrate do not change abruptly in the temperature range under consideration; abrupt changes might occur near a structural phase transition, for example. The variation in the specific heat and the thermal conductivity of the superconducting film can be ignored at the film thicknesses and laser pulse lengths involved here.

Figure 8 shows the factor by which the heating of the sample, ΔT_h, must exceed the heating observed experimentally, ΔT_{expt}, if the entire response of the system is to be attributed to the lattice heating. In the temperature range studied, however, we did not observe any abrupt change in the specific heat or thermal conductivity as the temperature was varied. This circumstance, combined with the temporal

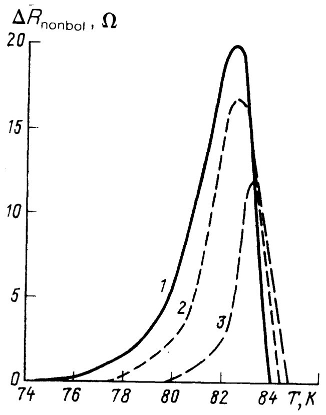

FIG. 7. Nonbolometric component of the photoresponse ΔR versus the sample temperature at various current densities j (A/cm^2). 1—2.6$\cdot 10^3$; 2—1.4$\cdot 10^3$; 3—3.1$\cdot 10^2$. The bridge thickness is h = 3300 Å.

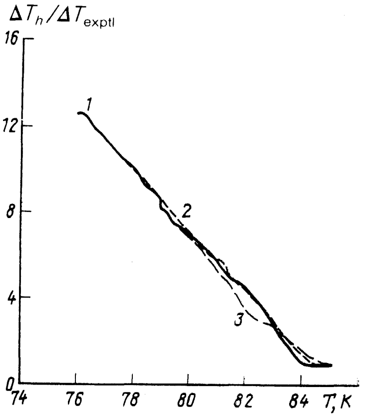

FIG. 8.

characteristics which we found, shows that the fast response of the system which we have singled out is of nonthermal origin.

The temperature dependence which we found for the amplitude of the photoresponse differs from the corresponding results of Refs. 8 and 9. Among the major differences are the decrease observed in Ref. 9 in the amplitude of the signal with increasing transport current and the broader temperature interval in which the nonthermal response exists.[8] In our own experiments we observed an increase in the amplitude of the signal with increasing transport current. We believe that these differences may stem from structural differences in the test samples, differences in film thickness, and differences in the dimensions of irregularities in the films. The experiments of Ref. 8 used films on MgO substrates. The characteristic size of the crystallites in superconducting films of that sort is usually larger than that in films on Zr_2O substrates. In addition, the film thicknesses in Ref. 8 were $d \sim 400$ Å, i.e., substantially less than those in the present experiments. The broadening of the temperature interval in which the nonbolometric response is observed with decreasing film thickness is confirmed by Fig. 9, which shows our results on the temperature dependence of the amplitude of the nonthermal response for two films differing in thickness.

The differences in the behavior of the amplitude of the response as a function of the transport current can be interpreted in the following way. Our experiments show that a threshold temperature T_1 is required for the appearance of a nonequilibrium resistive response of the system to optical excitation of a given intensity; i.e., no response is observed at $T < T_1$. If one assumes that the resistance which arises is associated with a breaking of weak links in the sample, then it is quite likely that the conditions prevailing in our experiments were such that the transport current promoted only a weakening of these links. As a result, the superconducting bridge would have become increasingly sensitive to the optical radiation with increasing transport current. At a high transport current[2)] the weak links may have been broken completely, with the result that the sensitivity of the bridge to the optical radiation fell off. These are probably the conditions which prevailed in the experiments of Ref. 8.

Analysis of the experimental results reveals another interesting aspect of the behavior of the nonthermal response. The nonthermal resistive response exists only when the bridge resistance $R(T)$ is below a certain R_0 (Fig. 10, a and b). At $R > R_0$, there is only a bolometric component. This effect is observed at various transport current densities. The value of R_0 depends on the film thickness and the transition width ΔT_c. For all the films studied, with various thicknesses, the value of $R_0/(R_N \Delta T_c)$, where R_N is the resistance of the sample in its normal (nonsuperconducting) state, is essentially a constant (within $\approx 5\%$). There is thus a relationship between R_0 and ΔT_c, i.e., the structure of the sample (the number and sizes of weak links and crystallites).

3. THEORY

We consider a model of a spatially nonuniform distribution of the order parameter $\Delta_0(x)$ in the superconductor (Fig. 11), and we assume that a "weak" superconductivity (in the regions between grains or blocks) is destroyed by even a weak external agent. Since the dimensions (a) of the crystallites are much larger than the distances (b) between them, the breaking of the weak links cannot lead to a large response of the system: The response will be small to the extent that the parameter b/a is. The breaking of the weak links in an inhomogeneous superconductor determines the

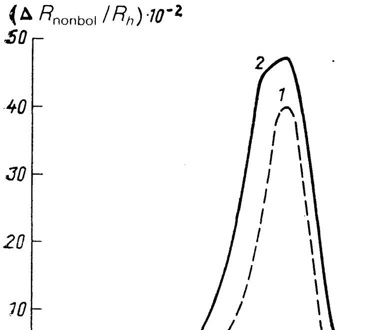

FIG. 9. Normalized amplitude of the nonbolometric signal versus the normalized temperature of the sample for various bridge thicknesses h (Å). 1—3300; 2—1500. The current density is $j = 1.4 \cdot 10^3$ A/cm^2.

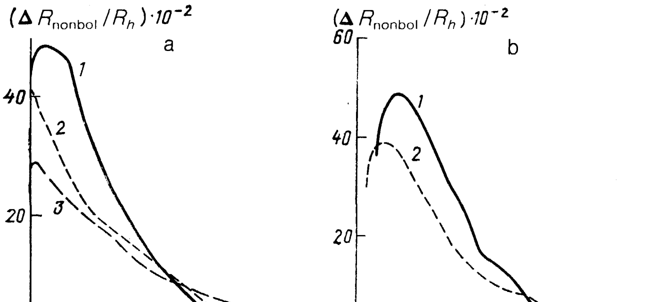

FIG. 10. Normalized nonbolometric response versus the resistance of the sample at various current densities j (A/cm^2). a: $h = 3300$ Å. 1—$2.6 \cdot 10^3$; 2—$1.5 \cdot 10^3$; 3—$3.4 \cdot 10^2$. b: $h = 6000$ Å. 1—$7.6 \cdot 10^3$; 2—$4.2 \cdot 10^3$.

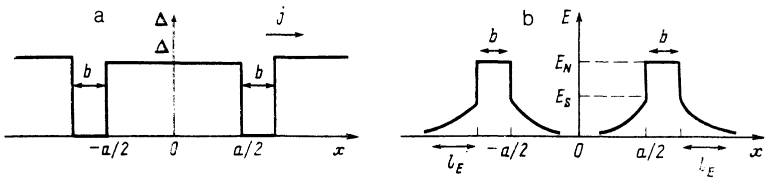

FIG. 11. a—Spatially inhomogeneous profile of the order parameter $\Delta(x)$; b—profile of the electric field in a spatially inhomogeneous superconductor having the profile $\Delta(x)$ shown in part a of this figure, as direct current flows through the superconductor.

threshold values of the external effects at which the resistive state arises in the superconductor. The destruction of superconductivity near the weak links may be caused by a coherence-instability mechanism.[3,5–7] Since Δ in these regions is small in comparison with the bulk value, the threshold light intensity for the appearance of the coherence instability may be low in them. Since the weak links are random, the threshold and the response of the system near the threshold are apparently determined by the statistics of these random links.

We are not interested in the behavior of the system near the threshold here. We are interested in the behavior of the resistive response at a high laser light intensity above the threshold. To study this behavior, we consider how the resistive response of the system with broken weak links changes, i.e., how the electric field penetrates into crystallites, $\sim a$ in size, in which the order parameter is quasiuniform. The penetration of an electric field into these regions may be associated with two effects: (1) a disbalance of the populations of the electron and hole branches of the quasiparticle spectrum[14–17] and (2) an Andreev reflection of quasiparticles from irregularities in the relief $\Delta(x)$ (Refs. 16–18). Each of these effects depends on the quasiparticle distribution function. By using external agents (e.g., optical radiation) to alter the quasiparticle distribution function, we can thus find different resistive responses of the system to this external agent.

a) Current-carrying states in inhomogeneous superconductors; Andreev reflection

Spatially inhomogeneous, current-carrying superconductors have been studied in connection with research on the intermediate states of type I superconductors.[13–15] The intermediate state has been represented as consisting of alternating layers of a superconducting phase S and a normal phase N. Our model of weak links in a superconductor is somewhat reminiscent of the structure of the intermediate state (there are alternating regions of weak and strong superconductivity), but there are differences: In our case the layers are not strictly periodic, the length scale of the regions of strong superconductivity satisfies $a \gg b$, where b is the characteristic size of the regions between the crystallites of the high T_c superconductor, and the value of the order parameter Δ in the S layers may not be small (as it is in the intermediate state). We accordingly use the model of an intermediate state which was proposed in Refs. 13–15, taking these differences into account, and also noting that the quasiparticle temperature in our case takes on nonequilibrium values and depends on the laser light intensity.

A kinetic equation for the quasiparticles was used in Refs. 13–15 to study the penetration of an electric field from N regions into a superconductor. The equation for the electrostatic potential is

$$\frac{d^2\Phi}{dx^2} - l_E^{-2}\Phi(x) = 0, \qquad (1)$$

where

$$l_E^2 = \begin{cases} L^2(T)\dfrac{4T}{\pi\Delta(T)}, & T \approx T_c \\ \dfrac{4\Theta_D^2\hbar v_F^2\tau}{3\pi\alpha_{ph}\Delta^3} \approx \mathrm{const}(T) + o\left(\dfrac{T}{\Delta}\right), & T \to 0, \end{cases}$$

T is the effective quasiparticle temperature, $\Delta(T)$ is the superconducting gap, $L(T)$ is the quasiparticle diffusion length, Θ_D is the Debye temperature, v_F is the Fermi velocity, $\alpha_{ph} \sim 1$, and τ^{-1} is the momentum relaxation frequency. At $T \approx T_c$ we have

$$L^2(T) = {}^1/_3 v_F^2 \tau\tau_\varepsilon(T), \quad \tau_\varepsilon^{-1} = 14\cdot\alpha_{ph}\zeta(3)\cdot T^3/\Theta_D^2,$$

where τ_ε^{-1} is the energy relaxation frequency.[16,17]

In the derivation of Eq. (1) in Refs. 15–17, it was assumed that the quasiparticles have a quasiequilibrium distribution function. It was also assumed that the parameter T in the length scale l_E is the temperature parameter in the Fermi distribution of quasiparticles. Generally speaking, a description of the quasiparticle distribution in the effective-temperature approximation is rather crude. In the case at hand, however, in which there is intense optical excitation (at an intensity above the threshold for the appearance of an nonbolometric response), and in which we are not interested in effects which stem from the details of the quasiparticle distribution function, that approach is valid.

The distance (l_E) which the electric field penetrates into the superconductor depends on Δ and also on the effective quasiparticle temperature T. Accordingly, it may vary if these parameters are varied (e.g., if the intensity of the optical excitation of the quasiparticles is varied). Estimating the length scale (l_E) to which the electric field penetrates into high T_c superconductors in Eq. (1) (we are taking the parameter values of $YBa_2Cu_3O_{7-x}$ from Refs. 11 and 19–22), we find $l_E \sim 100$ Å at $T \sim \Delta$. The length scale l_E may thus be less than or on the order of the length scales of the variations in the order parameter in the high T_c superconductors, $a \sim 100$–1000 Å. On this basis we can assume that the effects discussed in Refs. 13–15 may be important for the resistive photoresponse in these materials. Two parameters (and two effects) are associated with the effective quasiparticle temperature T, which depends on the light intensity, and with the value of the order parameter Δ: 1) the typical distance which the electric field penetrates into the superconductor, $l_E(T)$, and 2) the jump[15] in the electric field at the N–S layer boundary, which is a consequence of an Andreev reflection of quasiparticles from the nonuniform order parameter profile $\Delta(x)$.

Let us examine the spatially nonuniform order-parameter profile $\Delta(x)$ shown in Fig. 11a. The profile of the electric field which arises at such irregularities during the flow of an electric current is shown in Fig. 11b (according to Ref. 15). In accordance with (1), the electric-field profile in the interval $|x| < a/2$ is

$$E(x)=E_S \operatorname{ch}\left(\frac{x}{l_E}\right) \operatorname{ch}^{-1}\left[\frac{a}{2l_E}\right], \tag{2}$$

and the potential difference U_a on the interval from $x = -a/2$ to $x = a/2$ is found from

$$U_a=\int_{-a/2}^{a/2} E(x)\,dx=2l_E E_S \operatorname{th}\left(\frac{a}{2l_E}\right). \tag{3}$$

The potential difference over a region of the superconductor of length $a + b$ (a sort of average "period") is equal to simply the sum of the voltage drops across these regions:

$$U=U_a+U_b=2l_E E_S \operatorname{th}\left[a/(2l_E)\right]+bE_N. \tag{4}$$

From (4) we see that the length scale l_E appears in the expression for U. This length scale depends on the effective temperature, according to (1), and thus on the photoexcitation intensity. However, the Andreev reflection causes E_S to differ from the electric field in the normal (N) phase, $E_N = j/\sigma$, where j is the total current, and σ is the "metallic" conductivity, as was shown in Ref. 15. The jump $E_N - E_S$ in the electric field at the N–S boundary depends on the quasiparticle distribution function and thus on the effective temperature T. To find this dependence, we work from a kinetic equation for the quasiparticles. Solutions of this equation for the case in which there is an electric field in a superconductor in the intermediate state were studied in Refs. 13–15.

A solution of the kinetic equation for the quasiparticle distribution function was sought in Refs. 13–15 in the linear approximation:

$$n=n_F(\varepsilon)+\delta n\approx n_F(\varepsilon)+n_0(\xi, \mathbf{r})+n_1(\xi, \mathbf{r}) \cos\theta, \tag{5}$$

where θ is the angle in momentum space, reckoned from the current direction. Here

$$n_F(\varepsilon)={}^1/_2[1-\operatorname{th}(\varepsilon/2T)] \tag{6}$$

is a Fermi distribution with a renormalized spectrum of elementary excitations, and $\delta n/n \ll 1$ (Ref. 15, for example). Expression (5) is the standard Legendre-polynomial expansion of a nonequilibrium increment δn in a distribution function, in which we have retained only the first two terms. In this approximation the normal current j_n is expressed in terms of n_1 by

$$j_n(x)=e\int\frac{2\,d^3p}{(2\pi\hbar)^3}V_x n=\frac{eN}{p_F}\int d\xi\, n_1(\xi,x), \tag{7}$$

where the x axis has been chosen to run along the current direction, $V_x \approx V_F \cos\theta$, and $N = p_F^3/3\pi^2\hbar^3$. In the zeroth approximation in Δ/T, i.e., in the limit $\Delta \to 0$, which corresponds to a normal metal, the following expressions hold for an unbounded sample (Ref. 12, for example):

$$n_0(\xi)=-e\Phi\frac{\partial n_F}{\partial\varepsilon}\operatorname{sign}(\xi)=\frac{e\Phi}{4T}\operatorname{ch}^{-2}\left(\frac{\varepsilon}{2T}\right)\operatorname{sign}(\xi), \tag{8}$$

$$n_1(\xi)=-eEV_F\tau\frac{\partial n_F}{\partial\varepsilon}=\frac{eEV_F\tau}{4T}\operatorname{ch}^{-2}\left(\frac{\varepsilon}{2T}\right), \tag{9}$$

where at $\Delta = 0$ the quantity $\varepsilon = |\xi|$ is the quasiparticle energy in a normal metal.

Determining the functions n_0 and n_1 with the corresponding boundary conditions is an independent problem. It has been studied in particular in Refs. 14 and 15, in an effort to determine the distribution functions in the intermediate state in the S and N regions for above-barrier quasiparticles ($\varepsilon > \Delta$; see also Fig. 11a). Since an intermediate state was examined, and the parameter Δ/T was much less than unity, nearly all the quasiparticles were above-barrier particles ($\varepsilon > \Delta$). We know,[18] however, that the transmission coefficient for the movement of such particles out of an N region into an S region has a value $w < 1$. In other words, there is a finite probability that these particles will be reflected from the N–S boundary with changes in the sign of the momentum and the charge (this is Andreev reflection[18]). It was shown in Refs. 13–15 that Andreev reflection leads to a jump in the electric field at an N–S boundary (more precisely, it leads to an abrupt change in the electric field near the boundary, over a distance on the order of the coherence length). The physical meaning of the abrupt change in the electric field is that not all the quasiparticle current (the normal component of the current j_n) can flow across the boundary; part of it transforms into the superconducting current component j_s. The total current through a superconducting film, $j = j_n + j_s$, is thus a constant, but the relative values of j_n and j_s change at the N–S boundary.

The case in which we are interested here, of a spatially inhomogeneous superconductor, differs from the case of the intermediate state in a type I superconductor in that Δ has a highly nonuniform profile, of the type shown in Fig. 11a, even at $T \ll \Delta$. At $\Delta/T > 1$, the quasiparticle density in the weak-link regions (regions with a size on the order of b in Fig. 11a) is far higher than in regions of a strong superconductivity. As a result, the Andreev reflection at the N–S boundary is dominated not by the above-barrier quasiparticles ($\varepsilon > \Delta$; there are few quasiparticles at energies $\varepsilon > \Delta$) but by quasiparticles in the "wells" of the order-parameter profile $\Delta(x)$. As we have already seen, the mean free path in the high T_c superconductors satisfies $l_f \gtrsim \xi$, while the length scale of the variations satisfies $b \gg \xi$ [otherwise, there would be no wells on the $\Delta(x)$ profile, because of the proximity effect[12]]. The result is that in regions on the order of b in size the distribution function can be approximated in this case as being the same as in a bulk normal metal. We accordingly assume that Eqs. (8) and (9) hold for nonequilibrium increments in the distribution function.

In the N regions the superconducting current is zero ($j_s = 0$), so the total current is, according to (7),

$$j=j_n\approx\frac{eN}{p_F}\int_{-\infty}^{+\infty} n_1\,d\xi=\frac{2eN}{p_F}\int_0^{+\infty} n_1\,d\xi. \tag{10}$$

The normal component of the current, which stems from the particles in the well on the $\Delta(x)$ profile, transforms into a superconducting current at the boundary, since the transmission coefficient for these quasiparticles is $w = 0$ (Ref. 18). The magnitude of this current is

$$\tilde{j}_n{}^b = \frac{2eN}{p_F}\int_0^\Delta n_1\,d\xi \approx \sigma(E_N - E_S) \qquad (11)$$

(the total current in the N region is given by an equivalent expression, but the upper limit Δ in the integral would have to be replaced by $+\infty$). We have ignored the contribution to the Andreev reflection from above-barrier ($\varepsilon > \Delta$) quasiparticles under the assumption that the relation $\Delta/T \gg 1$ holds and under the assumption that there are relatively few such particles. As a result we find the following estimate of the jump in the electric field at the N–S boundary:

$$E_N - E_S \approx \frac{2eN}{\sigma p_F}\int_0^\Delta n_1\,d\xi = \frac{2}{eV_F\tau}\int_0^\Delta n_1(\xi)\,d\xi, \qquad (12)$$

where we have used $\sigma = Ne^2\tau/m$ for the conductivity, and we have taken $n_1(\xi)$ from (9).

Expression (12) is the same as the first term of Eq. (19) of Ref. 16. For the situation discussed in Ref. 16 (the intermediate state), that term played no major role and was accordingly discarded. In other words, the contribution of quasiparticles with energies $\varepsilon < \Delta$ to the Andreev reflection was discarded.

Substituting (9) into (12), we find (with $\Delta = 0$ and $\xi > 0$ we have $\xi \equiv \varepsilon$)

$$E_N - E_S \approx -2E_N\int_0^\Delta \frac{\partial n_F}{\partial\varepsilon}\,d\varepsilon = 2E_N[n_F(\varepsilon=0) - n_F(\varepsilon=\Delta)]$$

$$= E_N - 2E_N n_F(\varepsilon=\Delta). \qquad (13)$$

We thus find the following expression for the electric field in the superconductor, near the N–S boundary:

$$E_S = 2E_N n_F(\varepsilon=\Delta) = \frac{2E_N}{\exp(\Delta/T)+1}. \qquad (14)$$

It follows from (14) that at $\Delta/T \gg 1$ the electric field E_S is exponentially small: $E_S \approx 2E_N\exp(-\Delta/T)$.

Substituting (14) into (4), we find the potential difference across the region $a + b$:

$$U = E_N\left\{b + \frac{4l_E}{\exp(\Delta/T)+1}\,\mathrm{th}\left[\frac{a}{2l_E}\right]\right\}, \qquad (15)$$

where l_E is given by (1). Under the conditions $a/l_E \ll 1$ and $\Delta/T \ll 1$, expression (15) gives us (as it should) the potential difference across a film of a normal metal of length $a + b$:

$$U_N = E_N(b+a). \qquad (16)$$

b) "Hot" quasiparticles

Let us determine from the balance condition for the number of particles the extent to which the system of quasiparticles is heated by the laser light. We will make use of the results of Refs. 5–7, where solutions of the kinetic equation were analyzed for the case of optical excitation of quasiparticles, by a source described in the approximation $\omega \gg 2\Delta_0$ by

$$Q_\omega = \alpha_\omega\theta(\hbar\omega - \varepsilon), \qquad (17)$$

where

$$\alpha_\omega \approx \frac{2I}{c}\,\frac{e^2Lp_Fr}{m\hbar^2\omega^2(1+\omega^2\tau^2)}.$$

Here I is the energy flux density of the laser light; c is the velocity of light; $L(\frac{1}{3}v_F^2\tau\tau_{\mathrm{ph}})^{1/2}$ is the quasiparticle diffusion length; $v_F(p_F)$ is the velocity (momentum) of the electrons at the Fermi surface; $r = r_e r_{\mathrm{ph}}$; r_e is the quasiparticle breeding factor representing the breeding of quasiparticles by electron–electron collisions;[3] r_{ph} is the quasiparticle breeding factor for the breeding caused by the absorption of phonons; m is the effective mass of an electron; ω is the frequency of the optical excitation; τ is the momentum relaxation time for the relaxation caused by impurity centers.

We consider a quasihomogeneous superconductor with a superconducting gap Δ at $T < T_c$. We assume that the regions in Fig. 11a with the large value of the order parameter Δ (these regions have a typical size $\sim a$) are regions of a quasihomogeneous superconductor of precisely this sort. We assume that the quasiparticle energy distribution is

$$n_\varepsilon = n(\varepsilon) = [\exp(\varepsilon/T)+1]^{-1}, \qquad (18)$$

where T is the quasiparticle temperature, which is a function of the optical excitation intensity I. For the steady-state spatially uniform case the kinetic equation is[3]

$$(\partial n/\partial t)_{ph} + Q(\varepsilon) = 0, \qquad (19)$$

where the quasiparticle source $Q(\varepsilon)$ is given by (17). During optical excitation, the collision integral representing the collisions of quasiparticles with phonons is dominated by the recombination term S^R, as has been pointed out in Ref. 3, among other places. We also assume that the superconductor is at a low temperature, and we ignore the phonon distribution function, assuming $N_\varepsilon \ll 1$. We furthermore assume $k = -1$, where k is the power in the power-law dependence of the matrix element of the electron—phonon interaction on the wave vector, in the recombination term. We are making this assumption for simplicity; it was also made in Refs. 5–7. This assumption corresponds (Ref. 23, for example) to the case of PA phonons (piezoacoustic scattering). In this approximation we have the following expression for the recombination term:

$$S^R = \frac{1}{\tau_{ph}\Delta_0}\int_0^{\hbar\omega_D-\xi} d\xi'\left(1+\frac{\Delta^2}{\varepsilon\varepsilon'}\right)n(\varepsilon)n(\varepsilon'). \qquad (20)$$

In the approximation $T \ll \hbar\omega_D$ we can assume

$$\hbar\omega_D - \xi \approx \hbar\omega_D, \qquad (21)$$

since the quasiparticle distribution function is substantially nonzero at scales $\xi \sim T$ (Refs. 3 and 5–7). For the quasiparticle recombination rate R (the density of quasiparticles which recombine per unit time) we have thus

$$R = \frac{2m^2V_F}{\pi^2\hbar^3}\int_0^{\hbar\omega_D} S^R(\varepsilon)\,d\xi = \frac{2m^2V_F}{\pi^2\hbar^3\tau_{ph}\Delta_0}\int_0^{\hbar\omega_D} d\xi\, n(\varepsilon)$$

$$\times\int_0^{\hbar\omega_D-\xi} d\xi'\, n(\varepsilon')\left(1+\frac{\Delta^2}{\varepsilon\varepsilon'}\right)$$

$$\approx \frac{2m^2V_F}{\pi^2\hbar^3\tau_{ph}\Delta_0}\left\{\left(\frac{\pi^2\hbar^3N}{2m^2V_F}\right)^2 + \frac{\Delta^2}{4}\ln^2\left(\frac{\Delta}{\Delta_0}\right)\right\}, \qquad (22)$$

where, for PA phonons ($k = -1$),

$$\frac{1}{\tau_{ph}}=\frac{\pi\lambda}{2\hbar}\Delta_0, \tag{23}$$

$$N\equiv\frac{2}{(2\pi\hbar)^3}\int d^3p\, n(\varepsilon)=\frac{2m^2V_F}{\pi^2\hbar^3}\int_0^{\hbar\omega_D} n(\varepsilon)d\xi. \tag{24}$$

The second term in braces (curly brackets) in (22) stems from the one in the coherence factor $[1+(\Delta^2/\varepsilon\varepsilon')]$ in the collision integral. Physically, this term describes the change caused in the quasiparticle recombination rate by a deviation of the superconducting gap Δ from its equilibrium value Δ_0 at absolute zero. It can be seen from (22) that the coherence factor plays a particularly important role near the phase transition, where $\Delta\to 0$. At low quasiparticle temperatures, on the other hand, with $\Delta\approx\Delta_0$ and $\ln(\Delta/\Delta_0)\to 0$, the recombination rate is described by the usual expression (as, for example, in semiconductors) as a function of the quasiparticle density: $R\propto\gamma N^2$. It is important to note that γ may vary with the superconductivity mechanism in the various cases, while the dependence $R\propto N^2$ is universal. It reflects the fact that two quasiparticles participate in one recombination event.

Estimating the quasiparticle density, we finally find the following expression for the recombination rate:

$$R=\frac{\pi g v^2(\mu)}{2\hbar}\left\{\left[\frac{\Delta}{\exp(\Delta/T)+1}+T\ln\left(1+\exp\left(-\frac{\Delta}{T}\right)\right)\right]^2+\frac{\Delta^2}{4}\ln^2\left(\frac{\Delta}{\Delta_0}\right)\right\}, \tag{25}$$

where $v(\mu)=p_F m/\pi^2\hbar^3$ is the density of states at the Fermi surface and g is the matrix element of the electron–phonon interaction. For estimates we can assume[12] $gv(\mu)\sim 1$.

Integrating expression (17) for the quasiparticle source over the energy in a similar way, we find an average rate G of photoproduction of quasiparticles. This is the density of quasiparticles produced per unit time by the optical excitation:

$$G\approx\frac{2m^2v_F}{\pi^2\hbar^3}\int_0^\infty Q(\varepsilon)d\varepsilon\approx\frac{2m^2v_F}{\pi^2\hbar^3}\alpha_\omega\hbar\omega. \tag{26}$$

The intensity dependence of α_ω ($\alpha_\omega\propto I$) is taken from (17). Having expressions for the recombination rate (25) and the photoproduction rate (26) of the quasiparticles, we can then find from the condition for a steady state (the rates of photoproduction and recombination are equal),

$$R=G, \tag{27}$$

along with the equation for the superconducting gap Δ [see (6)], the dependence of the effective quasiparticle temperature T and the order parameter Δ on the light intensity I.

In writing the condition for a steady state in the form in (27) we are assuming that the quasiparticle density produced by the light is far higher than the equilibrium quasiparticle density. In general, however, we have to allow for the circumstance that the recombination rate in an equilibrium state is nonzero at a nonzero temperature, and on the left side of Eq. (27) we should have not the total recombination rate but its change from the equilibrium recombination rate R_t:

$$R-R_t=G, \tag{28}$$

where $R_t=R(T=T_t)$, T_t is the lattice temperature, and R is given by (25).

The case simplest to analyze is that in which the temperatures are low ($T_t\ll T_c$) and the light intensities are not too high. In this case the bulk value (i.e., the value inside the crystallites) of the order parameter is exponentially close to the value of the order parameter at a low temperature, Δ_0 (Ref. 12):

$$\Delta\approx\Delta_0-(2\pi\Delta_0T)^{1/2}\exp(-\Delta_0/T). \tag{29}$$

Since $\Delta\approx\Delta_0$, Eq. (17) by itself is sufficient for determining the effective temperature T in this case. We might add that the rate of "thermal" recombination is vanishingly low, $R_t\to 0$, and there is no need to generalize (28). In this case we have $T/\Delta\ll 1$; i.e., the bulk superconductivity is not destroyed. According to (25), the expression for the quasiparticle recombination rate is

$$R\approx\frac{\pi g v^2(\mu)}{2\hbar}\Delta_0^2\exp\left(-\frac{2\Delta_0}{T}\right). \tag{30}$$

Substituting (26) and (30) into (27), we find the following expression for the effective quasiparticle temperature:

$$T\approx 2\Delta_0\ln^{-1}\left[\frac{\Delta_0^2}{\hbar\omega\hbar\alpha_\omega}\right]=\frac{2\Delta_0}{\ln(I_0/I)}, \tag{31}$$

where the characteristic intensity is

$$I_0\approx\frac{cm\omega\Delta_0^2(1+\omega^2\tau^2)}{2e^2Lp_Fr}. \tag{32}$$

The weak (logarithmic) dependence of the effective temperature T on the light intensity I agrees well with the results of a numerical calculation,[24] where the phenomenological equation from Ref. 25 describing the quasiparticle diffusion was used, along with a corresponding diffusion equation for phonons.

An estimate of the value of I_0 for a neodymium laser (for light with a wavelength $\lambda_{las}=1.06\ \mu$m) and for the parameter values of the superconductor $YBa_2Cu_3O_{7-x}$ yields $I_0\sim 10^8$–10^9 W/cm^2. At reasonable light intensities we would thus have $I\ll I_0$. From (31), with photoexcitation intensities $I\sim 10^2$–10^3 W/cm^2, we find that the effective quasiparticle temperature T is on the order of 20–30 K. This is much lower than Δ_0. At $\Delta/T\gg 1$ we have

$$l_E^2=\frac{4\Theta_D^2\hbar v_F^2\tau}{3\pi\alpha_{ph}\Delta^3}\approx\text{const}. \tag{33}$$

According to (33), for $YBa_2Cu_3O_{7-x}$ (with the parameter values from Refs. 19–22, for example) we have $l_E\approx 20$–40 Å. Working from Eqs. (15) and (16), and taking account of our result in (31), we easily find the following expression for the photoresponse of the system under the conditions $\Delta/T\gg 1$ and $b/a\ll 1$:

$$\frac{U}{U_N}\approx\frac{1}{a+b}\left[b+4l_E\exp\left(-\frac{\Delta}{T}\right)\right]\approx\frac{b}{a}+\frac{4l_E}{a}\left(\frac{I-I_c}{I_0}\right)^{1/2},\quad I\geqslant I_c, \tag{34}$$

where I_c is the threshold intensity of the laser light. In (34) we have incorporated the circumstance that our estimates started from a state of the type shown in Fig. 11a, with broken weak links. It should thus be kept in mind that in our model there is a threshold value of the laser light intensity, I_c, such that the resistive response of the system arises only at $I > I_c$. This critical intensity I_c is determined by the statistics of the weak links and by the particular mechanism which breaks them.

It follows from (34) that near the threshold intensity I_c a voltage U arises abruptly: At $I = I_c$ we have $U/U_N \sim (b/a)/(1 + b/a)$. We can make use of this circumstance to estimate an average value of the parameter b/a from the experimental data. Figure 12 shows plots of $U(I)$ found from (34) for several parameter values. Expression (34) was derived for intensities $I \gtrsim I_c$.

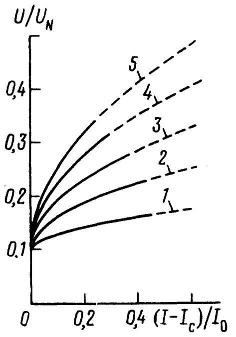

FIG. 12. Plots of expression (34) for various values of the parameter $4l_E/a$. 1—0.1; 2—0.2; 3—0.3; 4—0.4; 5—0.5. [$\Delta/T \gg 1$, $l_E \approx l_E$ $(T=0) = \mathrm{const}(T)$, $b/a = 0.1$].

The dependence of the photoresponse on the laser light intensity described by (34) (which was derived for the case $\Delta/T \gg 1$) stems from a change in the jump in the electric field at the N–S boundary upon a heating of the quasiparticles. This heating is determined by the Andreev reflection of quasiparticles at the N–S boundary. The length scale for the penetration of the field under the condition $\Delta/T \gg 1$ is $l_E = l_E(T) \approx \mathrm{const}(T) + o(T/\Delta)$.

As the effective temperature T increases, however, the change in the photoresponse voltage U begins to rise more rapidly, since the change in the field jump $E_N - E_S$ is now accompanied by a change in $l_E(T)$. In the T region close to T_c, the effect stems primarily from the change in the quantity

$$l_E(T) \sim [\Delta(T)]^{1/2} \sim [T_c - T]^{-1/4},$$

while the jump $E_N - E_S$ is close to zero, since the height of the barrier for the Andreev reflection approaches zero.

4. CONCLUSION

In summary, we have studied the properties of the nonequilibrium resistive response of thin films of the high T_c superconductor $YBa_2Cu_3O_{7-x}$ to external laser light. We have shown that this resistive response has components from both equilibrium (thermal) and nonequilibrium processes. Each of these processes is characterized by its own relaxation time. The thermal and nonthermal components of the response differ in the way they depend on the temperature, the light intensity, and the transport current flowing through the film. We have examined mechanisms for changes caused in the resistance of a spatially inhomogeneous superconducting film by a deviation from equilibrium in the quasiparticle system. We have shown that the change caused in the resistive state of an inhomogeneous superconductor by optical radiation may be a consequence of a penetration of an electric field into the superconductor. In addition to being of general physical importance (for research on the physics of nonequilibrium superconductivity in the high T_c materials), the effects of the hot quasiparticles may be of interest for applications, since the time scales for the establishment and destruction of the nonequilibrium state in the quasiparticle system can be very short ($\tau_\varepsilon \lesssim 10^{-12}$ s). This circumstance may be of assistance in the development of a variety of fast electronic devices using these effects.

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[1)] The superconductivity mechanism in the high T_c superconductors apparently might also be of a nonphonon nature, and not describable by the BCS model. However, (first) a self-consistent equation for the gap Δ corresponding to the BCS equation should be a part of any model. Second, in an analysis of highly nonequilibrium photoexcited states of a superconductor at $\hbar\omega \gg \Delta_0$ we are not particularly interested in the specific structure of the ground state or in the form of the equation for Δ.

[2)] One should of course bear in mind that the concepts of "large" and "small" transport currents are slightly arbitrary in this case and depend on the thickness and structure of the film.

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