Electron–phonon interaction in ultrathin Nb films

E. M. Gershenzon, M. E. Gershenzon, G. N. Gol'tsman, A. M. Lyul'kin, A. D. Semenov, and A. V. Sergeev

V. I. Lenin State Pedagogical Institute, Moscow

(Submitted 4 August 1989)
Zh. Eksp. Teor. Fiz. **97**, 901–911 (March 1990)

A study was made of the heating of electrons in normal resistive states of superconducting thin Nb films. The directly determined relaxation time of the resistance of a sample and the rise of the electron temperature were used to find the electron–phonon interaction time τ_{eph}. The dependence of τ_{eph} on the mean free path of electrons $\tau_{\text{eph}} \propto l^{-1}$ demonstrated, in agreement with the theoretical predictions, that the contribution of the inelastic scattering of electrons by impurities to the energy relaxation process decreased at low temperatures and the observed temperature dependence $\tau_{\text{eph}} \propto T^{-2}$ was due to a modification of the phonon spectrum in thin films.

1. Much new information on the electron–phonon interaction time τ_{eph} in thin films of normal metals and superconductors has been published recently.[1-7] This information has been obtained mainly as a result of two types of measurement. One includes experiments on weak electron localization investigated by the method of quantum interference corrections to the conductivity of disordered conductors,[1-3] which can be used to find the relaxation time τ_φ of the phase of the electron wave function. In the absence of the scattering of electrons by paramagnetic impurities the relaxation time τ_φ is associated with the most effective process of energy relaxation: $\tau_\varphi^{-1} = \tau_{\text{ee}}^{-1} + \tau_{\text{eph}}^{-1}$, where τ_{ee} is the electron–electron relaxation time. At low temperatures, when the dependence $\tau_{\text{ee}}^{-1} \propto T$ is exhibited by thin disordered films, the dominant channel is that of the electron–electron relaxation and there is a lower limit to the temperature range in which τ_{eph} can be investigated.

The low-temperature behavior of τ_{eph} is revealed by heating experiments in which the electron temperature increases by an amount $\Delta\theta \propto \tau_{\text{eph}}$ compared with the temperature T of phonons in thin films. The temperature range available to such experiments is limited from above by the phonon thermostat condition: the time for escape of a phonon τ_{es} from a film should be shorter than the phonon–electron scattering time τ_{phe}. In the case of films $d \sim 100$ Å thick if the phonon spectrum can be described by the Debye function, this condition imposes the limit

$$T < \hbar\eta v_F / k_B d \approx 10\ K$$

($\eta \sim 0.1$ is the coefficient representing the transmission of a phonon across the film–substrate interface and v_F is the Fermi velocity).

In recent investigations the electron temperature rise has been deduced from temperature-dependent quantum corrections to the resistance[4,5] and from the change in the resistance with temperature in the resistive state of a superconductor.[6,7] Since cooling of the electron subsystem is entirely due to the electron–phonon interaction, an important feature of such heating experiments is the ability to determine the value of τ_{eph} under conditions of strong electron scattering by impurities when the electron–phonon interaction is less effective than the electron–electron scattering, which then governs the value of τ_φ. It should be possible to use such experiments to identify the characteristic features of the electron–phonon interaction in "dirty" metals and in ultrathin films at low temperatures, when the condition $q_T l \ll 1$ is obeyed ($q_T = k_B T/\hbar u$ is the wave vector of a thermal phonon, u is the velocity of sound, and l is the mean free path of electrons).

The reported experiments[1-6] have revealed an interesting feature. In the case of ultrathin ($d \sim 100$ Å) metal films the temperature dependences $\tau_\varphi(T)$ and $\tau_{\text{eph}}(T)$ close to T^{-2} are usually observed in a wide temperature range $T = 1$–20 K. In the case of "localization" experiments such a dependence $\tau_\varphi(T)$ is usually and justifiably attributed to an intermediate asymptote of a more complex law $\tau_\varphi^{-1} = AT + BT^3$, where the first term corresponds to the electron–electron scattering in a thin disordered film and the second represents the electron–phonon scattering in the "pure" case characterized by $q_T l \gg 1$, but further investigations are needed to account for the dependences $\tau_{\text{eph}} \propto T^{-2}$ obtained from the heating experiments.

From the theoretical point of view, the low-temperature range is of particular interest. If the condition $ql \ll 1$ is satisfied, the electron–phonon interaction is modified significantly. In addition to the electron–phonon interaction processes, typical of a pure metal, we must take into account both the inelastic scattering of electrons by vibrating impurities, which is characterized by $\tau_{\text{eph}} \propto T^{-2} l$ (Ref. 8) as well as numerous electron–phonon interference processes. Consequently, the contribution of the inelastic scattering of electrons by impurities to the process of energy relaxation becomes smaller[9] and it is found that in the case of longitudinal and transverse phonons the temperature dependence is $\tau_{\text{eph}} \propto T^{-4} l^{-1}$. The same result is obtained in a reference system linked to the crystal lattice[10] in which case there is no inelastic scattering by impurities.

Different attempts to acount for the dependence $\tau_{\text{eph}} \propto T^{-2}$ at low temperatures have yielded different and inconsistent results:

a) if $ql \ll 1$, the contribution of the inelastic scattering of electrons by impurities ($\tau_{\text{eph}} \propto T^{-2}$) becomes smaller;

b) the view that the dependence $\tau_{\text{eph}} \propto T^{-2}$ corresponds to the pure electron–phonon interaction involving two-dimensional phonons in thin films is also inconsistent, because the condition $qd \ll 1$ that makes phonons two-dimensional leads naturally to the condition $ql \ll 1$ and in this case we have $\tau_{\text{eph}} \propto T^{-3} l^{-1}$ ($l < d$);

c) finally, the interaction of electrons due to the ex-

0038-5646/90/030505-07$03.00

change of virtual phonons does contribute to the dependence $\tau_\varphi \propto T^{-2}$ (Ref. 11), but it does not cool the electron subsystem and should not be manifested in the heating experiments.

In the pure case characterized by $ql \gg 1$ the interaction of electrons with transverse phonons is limited by the inelastic scattering of electrons on impurities ($\tau_{\rm eph} \propto T^{-2} l$) and can be ql times weaker than the interaction with longitudinal phonons ($\tau_{\rm eph} \propto T^{-3} l^0$). If the velocity of longitudinal sound u_l is much higher than the velocity of transverse sound u_t, then in a narrow range of temperatures where the longitudinal phonons satisfy the condition $q_l l \ll 1$, whereas the opposite condition $q_t l \gg 1$ is obeyed by transverse phonons, the contribution of the inelastic scattering by impurities to the energy relaxation process can be considerable.[12]

It must be stressed that these temperature dependences apply to the Debye phonon spectrum. At low temperatures the phonon spectrum of a thin film may be modified greatly and the degree of such modification depends on the acoustic matching of the film in question to its substrate, and also on the nature of defects and impurities present in the film. For example, in the case of a thin film ($dq_T \ll 1$) on an acoustically "soft" substrate the two-dimensional phonon spectrum is of the form $\omega_q \propto q^2$ (Ref. 13), which gives rise to the dependence $\tau_{\rm eph}^{-1} \propto T^{3/2} l (l \ll d)$. It is difficult to allow for the modification of the phonon spectrum *a priori* and to determine the range of the parameters $q(T)$ and l corresponding to the experimental conditions, i.e., whether we are dealing with the "dirty" limit, with the intermediate region, or with the "pure" case. Therefore, in addition to the temperature dependence of $\tau_{\rm eph}$ it is particularly desirable to study the dependence of $\tau_{\rm eph}$ on the mean free path of electrons l. Clearly, this dependence is governed primarily by the effective matrix element of the electron–phonon interaction: $|a|^2 \propto l^{-1}$ applies to the inelastic scattering of electrons by impurities and $|a|^2 \propto l$ corresponds to the "dirty" limit.

To the best of our knowledge the dependence on l has been investigated only for τ_φ. It was reported in Ref. 14 that in the temperature range 4.5–20 K the dependence of τ_φ on l was weak and it was found that $\tau_\varphi \propto T^{-2}$. This result was explained in Ref. 15 employing an intermediate asymptote and considering the case c) when the relaxation is due to the electron–phonon interaction.

The aim of the present investigation was to determine the dependence $\tau_{\rm eph}(l)$ at low temperatures using heating experiments. Electron heating was studied in superconducting Nb films so that a change in the resistance in the resistive state of the superconductor and the temperature-dependent quantum corrections to the resistance could both be used as an indicator of the electron temperature rise. This particular superconducting material was selected because it was possible to determine accurately the electron diffusion coefficient D from the temperature dependence of the second critical field. It should also be stressed that pure niobium has a complex Fermi surface, so that generally speaking we cannot apply the predictions of the theory of Refs. 9 and 10 to niobium. However, the investigated films had very low values of $l \sim 10$ Å and in this case significant disturbances of the spatial symmetry of the lattice resulted in isotropization of all the parameters of the material. On the other hand, in the metallic conduction range the smallness of l is not an obstacle in the interpretation of the experimental results on the basis of the theory of Refs. 9 and 10, because inclusion of the strong electron–impurity scattering (going beyond the Born approximation) gives the same dependences of $\tau_{\rm eph}$ on T and l. The influence of the substrate (via the phonon spectrum on the electron–phonon relaxation time was investigated by employing substrates with very different acoustic impedances: sapphire, silicon, quartz, glass, pyroceramic of the pyrosil type, and Mylar (polyethylene terephthalate).

2. We investigated Nb films formed by hf cathodic sputtering and electron-beam evaporation; they had various thicknesses and different values of the mean free path of electrons. In the sputtering and electron-beam evaporation methods we used the same starting material which was Nb of 99.9999% purity with the resistance ratio $R(300)/R(10) = 20\,000$. Our samples were strips $W = 1$ μm wide and $L = 100\,\mu$m long formed by photolithography. The thickness d of the films prepared by hf sputtering was determined from the sputtering rate and duration; in the electron-beam evaporation case we determined d using a quartz thickness meter. The parameters of some of the investigated films are listed in Table I (T_c is the temperature of the superconducting transition and $R_\square$ is the resistance per square of the film).

The conduction process was strongly metallic and the temperature dependence of the conductivity was described well by the weak localization theory.

The electron diffusion coefficient D was determined for these films from the temperature dependence of the second critical field H_{c2} using the expression

$$D=1{,}27\frac{ck_B}{e}\left[\left(\frac{dH_{c2}}{dT}\right)_{T=T_c}\right]^{-1}. \qquad (1)$$

It should be pointed out that the popular methods for the determination of l contain empirical constants obtained for much purer niobium and they describe poorly our films. In the case of relatively pure Nb the product ρl increases on reduction in l, indicating a reduction in the density of states ν_0 of electrons on the Fermi surface.[16] In the other method, based on calculation of the resistance drop as a result of cooling from room to helium temperatures, use is made of the mean free path of electrons governed by the electron–phonon scattering, which again depends strongly on the electron parameters ν_0 and v_F.

In the case of our films the electrical resistivity obeyed $\rho \propto D^{-1}$, as demonstrated in Fig. 1, and $\nu_0 = (e^2 \rho D)^{-1} = 1.6\times 10^{23}$ eV$^{-1}\cdot$cm^{-3} was independent of D. Consequently, we assumed that the change in the structure of the Fermi surface occurred when the defect concentration was much less than in our films and the electron parameters ν_0 and v_F were not affected by further reduction in l. This density of states corresponded to the Sommerfeld constant $\gamma = (6.2 \pm 0.4)\times 10^{-4}$ J$\cdot$cm$^3\cdot$K^1, which differed little from $\gamma = 6.65\times 10^{-4}$ J$\cdot$cm$^3\cdot$K^1 for pure niobium. In the case of films characterized by the mean free path $l \sim 100$ Å the change in the electron parameters was large ($\sim 10\%$).[16]

Clearly, the nonmonotonic variation of the density of the electron states on reduction in l was due to the following reasons. If $l \gg 100$ Å, the electron–impurity scattering smears out the density-of-states peak created at the Fermi level of pure niobium and in the case of our films the energy interval

TABLE I.

No. п/п	d, Å	D, cm^2/s	$R_\square$, Ω	T_c, К	Substrate
1	120	0.75	44	3,2	sapphire
2	80	0.80	52	4,9	sapphire
2a	80	0.80	55	4,5	silicon
3	150	1.30	18	5,9	sapphire
4	180	0.72	25	4,5	glass
5	180	0.70	28	4,3	Mylar
6	50	1,2	50	4,4	sapphire
7	200	0,90	23	5,5	sapphire
8	400	1,65	6,0	7,0	sapphire
9	600	1,0	6,8	6,7	sapphire
10	1000	2,65	1,5	8,5	sapphire

$\Delta\varepsilon \sim \hbar v_F/l \sim 1$ eV includes already the adjacent density-of-states peaks. In view of the constancy of v_F, we shall assume that the variable in the experimental dependences is the diffusion coefficient D and in discussions we shall use the mean free path $l = 3D/v_F$.

We determined the electron–phonon interaction time $\tau_{\rm eph}$ using both methods mentioned in Sec. 1. In the localization experiments the time $\tau_{\rm eph}$ was calculated from the magnetic-field dependences of the quantum corrections to the conductivity; in the case of the heating experiments we used directly the frequency dependences of the response of a sample to low-intensity electromagnetic radiation and made calculations employing the value of $\Delta\theta$ obtained under steady-state conditions. We shall not discuss details of the first method, dealt fully in the literature (see, for example, Refs. 2 and 3), but we shall consider only the second method.

In the case of sufficiently thin films at low temperatures the escape time of phonons from a film $\tau_{\rm es} = 4d/\eta u$ is less than the phonon–electron scattering time $\tau_{\rm phe}$. Then, the nonequilibrium phonons emitted by electrons leave the sample and have no influence on the electron subsystem, so that the phonons in the film act as a thermostat. Moreover, under the conditions of strong impurity scattering at sufficiently low temperatures the electron–electron interaction is more effective than the electron–phonon interaction and a nonquilibrium distribution function of electrons established by electromagnetic radiation is of the Fermi nature with a temperature θ exceeding the thermostat temperature T (Ref. 17). Then, if $\theta - T \ll T$ the temperature of the electron subsystem exposed to electromagnetic radiation is governed by the heat balance equation

$$c_e\frac{d\theta}{dt}=\rho(\theta,j)j^2-\frac{c_e}{\tau_\theta}(\theta-T)+P_\sim, \qquad (2)$$

where c_e is the electron specific heat; j is the density of the transport current; $P_\sim = P_0(1+\cos 2\pi ft)$ is the power of the electromagnetic radiation absorbed per unit volume of the film; f is the amplitude modulation frequency; τ_θ is the electron temperature relaxation time representing the electron–phonon interaction time averaged over the energy ε and differing from $\tau_{\rm eph}$ for electrons on the Fermi surface ($\varepsilon = 0$) by a numerical factor smaller than 1. In the future we shall not distinguish between τ_θ and $\tau_{\rm eph}$.

In this case the amplitude of the component of θ alternating at the modulation frequency is given by

$$\Delta\theta=\frac{P_0\tau_{eph}}{c_e}[1+(2\pi\tau_{eph}f)^2]^{-1/2}. \qquad (3)$$

Since the resistance of a sample is a function of temperature, the interaction with electromagnetic radiation at a constant value of the current I alters the voltage across it $\Delta U = \Delta\theta\cdot dU/dT$ ($dU/dT = I\cdot dR/dT$ is the temperature coefficient of the voltage).

In our experiments we used radiation emitted within a narrow range of millimeter and submillimeter range of wavelengths. Amplitude modulation at frequencies $f = 10^4$–10^9 Hz was due to beats of the oscillations generated by two backward-wave tubes (the frequency shift of one of them was controlled). The phase self-matching system for the beat frequency ensured that the frequency was stable to within better than 10^{-8} and, therefore, it was possible to reduce the width of the measurement band to 10 Hz throughout the investigated range of f (Ref. 18). This method made it possible to determine directly the relaxation time and it was much more sensitive than the widely used pulse technique. In the latter case the radiation sources not only emit short pulses, but these pulses are characterized by a large off-duty factor, so that it is necessary to record the response in a fairly wide frequency band. This means that the radiation has to be of sufficiently high power and the net result is overheating of the sample and nonlinear relaxation of the electron temperature.

Below the critical superconducting temperature we ensured sufficiently high values of dU/dT by making direct measurements of $\tau_{\rm eph}$ in the resistive state of a superconductor, which was induced by the application of a magnetic field. We demonstrated that the specific nature of the resis-

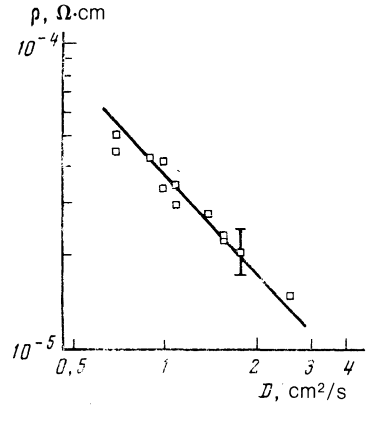

FIG. 1. Dependence of the electrical resistivity ρ on the diffusion coefficient D of electrons in thin ($d = 50$–1000 Å) Nb films. The continuous line represents the dependence $\rho \propto D^{-1}$.

tive state of the superconductor had practically no influence on the energy relaxation rate by making a number of measurements using samples converted completely to the normal state by the applied magnetic field.

In addition to the direct determination of $\tau_{\rm eph}$ we used the dependence $\Delta U(f)$ to find its value by a quasisteady-state method involving measurement of dU/dT and ΔU at a frequency $f \ll (2\pi\tau_{\rm eph})^{-1}$. Since $c_e = \gamma T$, we found that

$$\tau_{eph} = \frac{\Delta U \gamma T}{(dU/dT) P_0}. \tag{4}$$

This quasisteady-state method was more sensitive and it enabled us to use samples exhibiting an extremely weak temperature dependence of the resistance $[(dR/dT)R^{-1} \approx 10^{-4} K^{-1})]$.

In the determination of $\tau_{\rm eph}$ by the quasisteady-state method it was important to determine the radiation power P_0 absorbed per unit volume of the film. This was done in the infrared range in free space, as well as in the 2-mm range of wavelengths when the investigated sample was inside a waveguide. In the former case the source was an absolute black body and the power of the radiation reaching a sample was found from geometric considerations; in the latter case the radiation power was measured using a thermistor bridge. The absorption coefficient of the investigated sample (which was a structure consisting of a thin metal film) was calculated in both cases allowing for the diffraction.[19,20] Another method for the determination of the absorption coefficient used in the present study involved a comparison of P_0 with the power dissipated by the current. In our experiments we ensured the same change in P_0 under the influence of the hf and static fields: $\Delta\tau_{\rm eph}(P_0) = \Delta\tau_{\rm eph}(\rho j^2)$, where $\tau_{\rm eph}$ was determined by a transient method and the value of the resistivity ρ was kept constant and equal to the value at low currents by varying the magnetic field.[21] The results of our determination of P_0 by these two methods agreed to within 20% if in the second method the measurements were carried out under conditions such that the resistance of the film was no more than twice the resistance in the normal state.

These experiments enabled us also to select in a sound manner the values of P_0 and j so that the electron temperature rise $\Delta\theta = \theta - T$ was sufficiently small and did not distort the temperature dependence of $\tau_{\rm eph}$. With the exception of the experiments involving determination of the absorption coefficient the measurements carried out in the temperature range from 1.5 to 10 K were carried out under such conditions that the densities of the Joule power and of the power of the adsorbed electromagnetic radiation did not exceed 10^{-4}–3×10^{6} W/cm^3, which ensured that $\Delta\theta \leqslant 0.1$ K.

Two methods were used to check that the phonon thermostat condition was satisfied. In the first method the dependences $\Delta U(f)$ were determined by varying the heat transfer conditions (transparency η): the same sample was studied in vacuum or in liquid helium, and similar samples on different substrates were used. In the second method we used the dependences $\tau(d)$ to find directly the range of film thicknesses d where the condition $\tau_{\rm es} \ll \tau_{\rm eph}$ was obeyed.

3. The amplitude-frequency dependences of the response of several similar Nb films on different substrates were obtained at $T = 1.5$ K (Fig. 2). The samples were kept in vacuum and heat was transferred from the substrate via massive copper blocks. At high frequencies the $\Delta U(f)$ curves obeyed

$$\Delta U = \Delta U_0 [1 + (2\pi f \tau)^2]^{-1/2}$$

[see Eq. (3)] and were practically identical for samples on different substrates throughout the investigated temperature range. The values of $f_0 = 1/2\pi\tau$ are identified in Fig. 2 by arrows. This gave the values of τ for thin films, which were identified by us with the electron–phonon interaction time $\tau_{\rm eph}$.

The monotonic fall of ΔU on increase in the frequency at low values of f was due to the reverse flux of nonequilibrium phonons from the substrate to the film, i.e., it was governed by the thermal conductivity of the substrate. In the case of films deposited on sapphire substrates the effect was small and did not appear even at modulation frequencies in the range $f > 100$ Hz. When a sample was immersed in liquid helium the lf "tails" disappeared irrespective of the substrate material, whereas the hf parts of the dependences $\Delta U(f)$ were unaffected. Consequently, in measurements by the quasisteady-state method the frequency f of modulation of the incident radiation was always selected to be sufficiently high so as to avoid the influence of the reverse flux of nonequilibrium phonons from the substrate. The powers used in our experiments were such that the overall increase in the temperature of the film relative to the thermostat did not exceed 0.1 K.

The temperature dependences of the time τ, identified by us as $\tau_{\rm eph}$, are plotted for two Nb samples in Fig. 3. The continuous curves correspond to the dependence $\tau_{\rm eph} \propto T^{-2}$. It should be pointed out that the results obtained by the quasisteady-state method agreed with the directly measured values of $\tau_{\rm eph}$ in the resistive state, although they were obtained in a wider range of temperatures (up to 10 K). The values of $\tau_{\rm eph}$ determined at $T = 1.6$–2 K from the dependence $\Delta U(f)$ in the normal state also agreed with the values of $\tau_{\rm eph}$ found in the resistive state using fields in the range $H = (0.8$–$1])H_{c2}$. The discrepancy of τ near the critical temperature obeyed a dependence $\tau \propto (1 - T/T_c)^{-1/2}$ (inset in Fig. 3), typical of the relaxation time τ_Δ of the order parameter. We included in this figure also the values of τ_φ deduced from the magnetic-field dependences of the quantum corrections to the conductivity. The dashed straight-lines represent the dependence $\tau_\varphi \propto T^{-1}$.

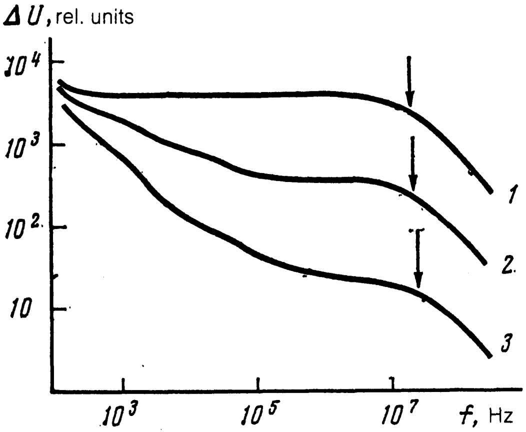

FIG. 2. Dependences $\Delta U(f)$ obtained for Nb samples on different substrates (see Table I): 1) No. 2, No. 2a; 2) No. 4; 3) No. 5; $T = 1.5$ K.

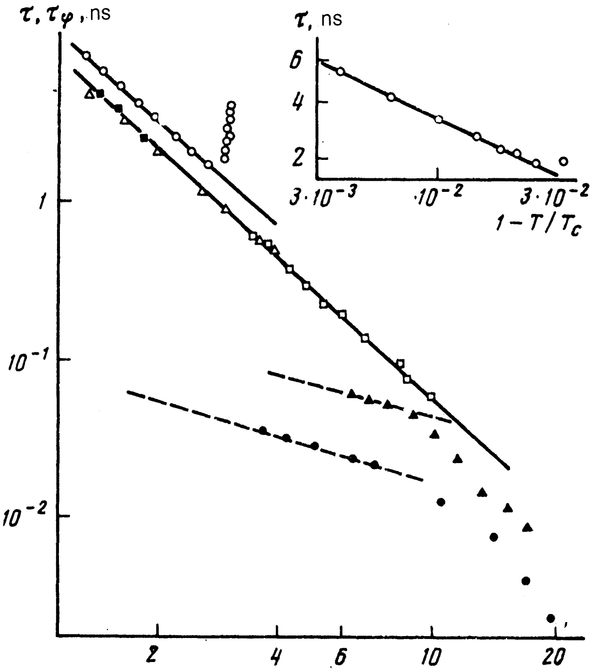

FIG. 3. Temperature dependences of τ (○, △, ■, □) and τ_φ (●, ▲) obtained for samples No. 1 (○, ●) and No. 3 (△, ▲, ■, □). The results were deduced from the measured dependences $\Delta U(f)$ obtained in the resistive state (○, △), in the normal state when $H > H_{c2}(T)$ (■), and from the temperature dependences of ΔU and dU/dT in the normal state (□). The continuous curves are the dependences $\tau \propto T^{-2}$, whereas the dashed curves represent $\tau \propto T^{-1}$. The inset shows the dependence $\tau(T)$ near T_c for sample No. 1 [the continuous curve represents the dependence $\tau \propto (1 - T/T_c)^{-1/2}$].

The results plotted in Fig. 3 demonstrate that measurements based on the electron heating gave the dependence $\tau_{\text{eph}} \propto T^{-2}$ in a wide temperature range $T = 1.5$–10 K, whereas at temperatures $T < 10$ K we found that $\tau_\varphi \propto T^{-1}$, corresponding to the electron–electron scattering. At higher temperatures the ratio of the times τ_{ee} and τ_{eph} varied smoothly and at $T > 12$ K the rate of relaxation of the wave-function phase, governed by the fastest inelastic relaxation process, was determined primarily by the electron–phonon interaction. The values of τ_φ obtained at $T > 12$ K were in approximate agreement with the values of τ_{eph} extrapolated to this range and deduced from the heating experiments.

Figure 4 shows the dependence $\tau(D)$ obtained at $T = 1.6$ K; the thickness of the samples used in this case did not exceed 200 Å. The results of the measurements indicated that the dependence $\tau_{\text{eph}} \propto D^{-1}$ was obeyed by these films.

The results of an investigation of $\tau(T)$ for films of different thickness d are plotted in Fig. 5. An analysis of these results demonstrated that the temperature dependence $\tau(T)$ became modified beginning from a certain thickness. In the case of films with $d = 50$–200 Å the cooling time τ was described by the expression $\tau \propto T^{-p}$, where $p = 2.0 \pm 0.2$, whereas in the case of the thicker samples the dependence $\tau(T)$ could no longer be described by a power law. However, if it was still approximated by the same method as before, it was found that the power exponent decreased smoothly and for a film of thickness $d = 1000$ Å its average value was close to 1.

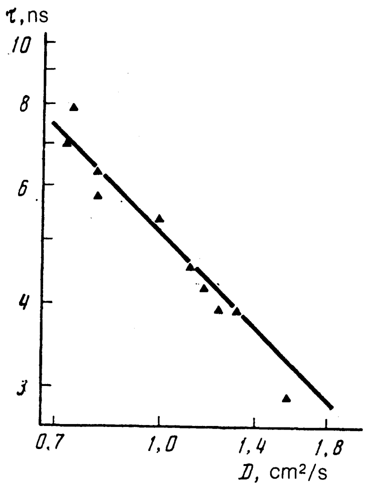

FIG. 4. Dependence $\tau(D)$ determined at $T = 1.6$ K for samples of thickness $d \leqslant 200$ Å; the continuous curve represents the dependence $\tau \propto D^{-1}$.

In the case of the bolometric effect, expected for the thickest films, we should have

$$\tau = \tau_b = (C_e/C_p)\, \tau_{\text{es}} \propto d$$

(τ_b is the cooling time of the film as a whole, i.e., the bolometric time). In the case of thin films we observed $\tau_{\text{eph}} \propto D^{-1}$ (Fig. 4); in the transition region the value of $\tau = \tau_b + \tau_{\text{eph}}$ depended on both D and d. We separated the dependence of τ on d by plotting the results in Fig. 6 in such a way as to reduce τ to the same diffusion coefficient $D = 1.0$ cm²/s⁻¹. It is clear from Fig. 6 that the thickness d_0 corresponding to the condition $\tau_b = \tau_{\text{eph}}$, deduced from the intersection of the straight lines $\tau \propto d$ and $\tau = \tau_{\text{eph}}$ ($D = 1.0$ cm²/s) obtained from the experimental data, increased as a result of cooling ($d_0 = 360$ Å at $T = 4.2$ K and $d_0 = 680$ Å at $T = 1.6$ K). Using the ratio $C_e/\tau_{\text{eph}} = C_{\text{ph}}/\tau_{\text{phe}}$, we could rewrite the condition $\tau_b(d_0) = \tau_{\text{eph}}$ in the form

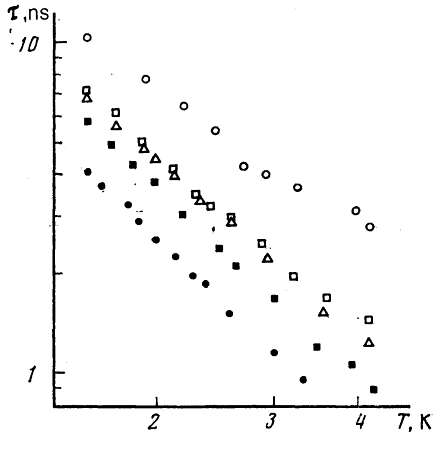

FIG. 5. Temperature dependences of τ obtained for samples of different thickness and with different diffusion coefficients: ●) sample No. 6; ■) No. 7; △) No. 8; □) No. 9; ○) No. 10.

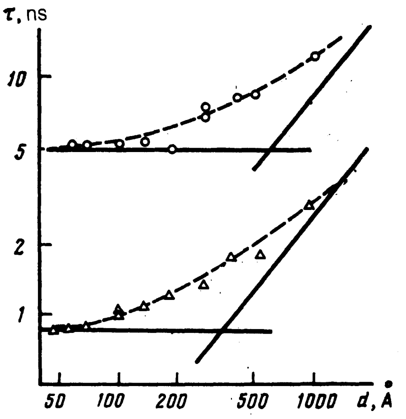

FIG. 6. Dependences $\tau(d)$ obtained for $D = 1.0\ \mathrm{cm}^2/\mathrm{s}^{-1}$ at two temperatures: ○) 1.6 K; △) 4.2 K. The dashed curves are plotted by extrapolation of the experimental data using the appropriate dependence of $\tau = \tau_{eph} + \tau_b$; the continuous lines represent the dependences $\tau = \tau_b \propto d$ and $\tau = \tau_{eph}$.

$\tau_{es}(d_0) = \tau_{phe}(T)$, which made it possible to estimate the phonon–electron scattering time for the investigated films and the temperature dependence of this time.

4. As pointed out already, the heating effect was deduced in our experiments not only from the change in the resistance in the resistive state of the investigated state of the superconductor, but also from the temperature-dependent quantum corrections to the resistance. In the former case the steep slope of the temperature dependence of the resistance made it possible to detect readily even small changes in the electron temperature. In the latter case the slope was many orders of magnitude less so that we had to allow also for other possible effects apart from the electron heating.

Suppression of the localized contribution to the conductivity $\Delta\sigma_{loc}$ by hf heating did not affect the results of measurements of τ_{eph} in the normal state and this was due to a number of reasons. It is sufficient to point out that in the range $T_c < T < 3T_c \approx 12$ K the measurements were carried out when $B = 0$ and also in a magnetic field ($B = 5\ T$) sufficiently strong to suppress the superconducting fluctuations and also to suppress $\Delta\sigma_{loc}$. It was found that the results of measurements of τ_{eph} were the same for $B = 0$ and $B = 5\ T$.

Another possible reason for the magnetic-field dependence of the resistance, unrelated to heating, is a change in the intergranular tunneling. However, it is known to occur in thin films in electric fields of much higher intensity[22] than those used in our experiments. Moreover, the characteristic frequencies at which the dispersion of this contribution to the resistance is observed exceeds considerably the values of τ_{eph}^{-1} obtained in the present study. Therefore, the change in the resistance is a consequence of the electron heating also in the normal state.

These experimental results demonstrate that all the above electron heating conditions had been satisfied. In fact, it follows from the data in Fig. 6 that if $d < d_0$ the values of τ no longer depend on d, i.e., τ_{es} becomes less than τ_{phe} We can use Fig. 3 to obtain the values of τ_{ee} (dashed lines) and confirm that in the range $T < 10$ K the inequality $\tau_{ee} \ll \tau_{eph}$ is obeyed by our samples, so that the nonequilibrium electron distribution function is of the Fermi nature and is characterized by an electron temperature θ.

The values of the critical thickness $d_0(T)$, corresponding to the transition from the bolometric effect to the electron heating, i.e., $\tau_{es}(d_0) = \tau_{phe}(T)$, were found from the experimental data plotted in Fig. 6 and allowed us to estimate the temperature dependence of the phonon–electron scattering time

$$\tau_{phe}(T) = 4d_0(T)/\eta u \propto T^{-0.7}.$$

The main result of the present study is the dependence of the energy relaxation of electrons interacting with phonons on the temperature and diffusion coefficient of electrons deduced from the electron heating in thin Nb films: $\tau_{eph} \propto T^{-2}D^{-1}$ in the range $1.5\ \mathrm{K} < T < 10$ K and $0.6 < D < 1.8$ cm^2/s. This dependence was deduced from directly determined time dependences and from the electron heating under steady-state conditions. We were of the opinion that the change in the diffusion coefficient and in the time τ_{eph} is primarily due to the change in the mean free path of electrons l, whereas the characteristic of the electron spectrum deduced from the dependence $\rho(D)$ (Fig. 1) are practically unaffected in the range of values of l typical of our samples. Then, the dependence $\tau_{eph} \propto l^{-1}$ agrees with the theoretical dependence of the effective matrix element of the electron–phonon interaction on condition that $q_T l \ll 1$ and the temperature dependence $\tau_{eph} \propto T^{-2}$ is due to the modification of the phonon spectrum (for example, that resulting from the contribution of flexural vibrations-characterized by the dispersion law $\omega \propto q^2$). In fact, an estimate of the phonon specific heat deduced from our data

$$c_{ph} \approx c_e \tau_{phe}/\tau_{eph} \propto T^2$$

shows that the phonon spectrum of a film differs considerably from the Debye form. This interpretation is supported also by the results of other experiments of the heating type at low temperatures: the power exponent p in the temperature dependence of τ_{eph}^{-1}, close to 2, has been reported for ultrathin Au films of thicknesses $d = 30$ Å (Ref. 5) and $d = 75$ Å (Ref. 4), whereas in the case of AuPd films of $d = 310$ Å the power exponent is $p \approx 3.7$ (Ref. 23). On the other hand, it follows from our experiments that the value of τ_{eph} is independent of the substrate material and this is due to the fact that the modification of the phonon spectrum in Nb films is governed by the film structure itself and is independent of the conditions of the acoustic matching at the film–substrate interface. Investigations of the phonon spectra of thin films are needed in order to obtain specific information on this topic.

The results obtained demonstrate that the temperature dependence $\tau_{eph} \propto T^{-2}$ observed in many experiments is not related to the inelastic scattering of electrons by impurities (for then we should have $\tau_{eph} \propto l$) and, in agreement with the theoretical predictions; the contribution of such processes to the energy relaxation at low temperatures becomes less compared with the contributions of other electron–phonon–impurity interference processes.[9,10]

We shall conclude by noting that an investigation of τ_{eph} on the basis of the electron heating in the resistive state of superconducting films is interesting also from the practical point of view. This is because the investigated samples are used as sensitive elements in superconducting bolometers with a response time close to the theoretical limit.[24] A direct determination of τ_{eph} and its dependences on the parameters of the material and on the external conditions should make it possible to predict the characteristics of such bolometric devices.

This investigation was supported by the Scientific Council on the Problem of High-Temperature Superconductivity and was carried out as part of the project No. 40 of the "High-Temperature Superconductivity" State Program.

[1] G. Bergmann, Phys. Rep. **107**, 1 (1984).
[2] Proc. Intern. Conf. on Localization, Interaction and Transport Phenomena in Impure Metals, Brunswick, West Germany, 1984 (ed. by B. Kramer, G. Bergmann, and Y. Bruwnseraede), Springer Verlag, Berlin (1985) [Springer Series in Solid-State Sciences, Vol. 61].
[3] Yu. V. Sharvin and D. Yu. Sharvin, in *Low Temperature Physics* (ed. by A. S. Borovik-Romanov), Mir Publishing Press, Moscow (1985), p. 221.
[4] S. I. Dorozhkin, F. Lell, and W. Schoepe, Solid State Commun. **60**, 245 (1986).
[5] G. Bergmann, Solid State Commun. **46**, 347 (1983).
[6] E. M. Gershenzon, M. E. Gershenzon, G. N. Gol'tsman, A. D. Semenov, and A. V. Sergeev, Zh. Eksp. Teor. Fiz. **86**, 758 (1984) [Sov. Phys. JETP **59**, 442 (1984)].
[7] J. P. Maneval, J. Desailly, and B. Pannetier, Proc. Fourth Intern. Conf. on Phonon Scattering in Condensed Matter, Stuttgart, West Germany, 1983, publ. by Springer Verlag, Berlin (1984), p. 43.
[8] H. Takayama, Z. Phys. **263**, 329 (1973).
[9] M. Yu. Reĭzer and A. V. Sergeev, Zh. Eksp. Teor. Fiz. **90**, 1056 (1986) [Sov. Phys. JETP **63**, 616 (1986)].
[10] A. Schmid, Z. Phys. **259**, 421 (1973).
[11] D. Belitz, Phys. Rev. B **36**, 2513 (1987).
[12] J. Rammer and A. Schmid Phys. Rev. B **34**, 1352 (1986).
[13] I. M. Lifshitz, Zh. Eksp. Teor. Fiz. **22**, 475 (1952).
[14] R. P. Peters and G. Bergmann, J. Phys. Soc. Jpn. **54**, 3478 (1985).
[15] D. Belitz, S. Das Sarma, Phys. Rev. B **36**, 7701 (1987).
[16] V. M. Pan, V. G. Prokhorov, and G. G. Kaminskiĭ, Fiz. Nizk. Temp. **6**, 968 (1980) [Sov. J. Low Temp. Phys. **6**, 470 (1980)].
[17] E. M. Gershenzon, M. E. Gershenzon, G. N. Gol'tsman, A. D. Semenov, and A. V. Sergeev, Pis'ma Zh. Eksp. Teor. Fiz. **36**, 241 (1982) [JETP Lett. **36**, 296 (1982)].
[18] E. M. Gershenzon, G. N. Gol'tsman, and G. I. Mirskiĭ, Prib. Tekh. Eksp. No. 4, 131 (1987).
[19] I. M. Starobin, Radiotekh. Elektron. No. 1, 80 (1990).
[20] I. M. Starobin, Radiotekh. Elektron. No. 2, 141 (1990).
[21] E. M. Gershenzon, G. N. Gol'tsman, A. I. Elant'ev, *et al.*, Fiz. Nizk. Temp. **14**, 753 (1988) [Sov. J. Low Temp. Phys. **14**, 414 (1988)].
[22] M. E. Gershenzon and V. N. Gubankov, Pis'ma Zh. Eksp. Teor. Fiz. **34**, 32 (1981) [JETP Lett. **34**, 30 (1981)].
[23] R. A. Lee and M. N. Wybourne, J. Phys. F **16**, L169 (1986).
[24] E. M. Gershenzon, M. E. Gershenzon, G. N. Gol'tsman, *et al.*, Zh. Tekh. Fiz. **59**(2), 111 (1989) [Sov. Phys. Tech. Phys. **34**, 195 (1989)].

Translated by A. Tybulewicz