Influence of optical pumping on the tunnel characteristics of superconductors

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It is shown that under the influence of optical pumping the peak in the phonon density of states of a superconductor at frequency V_0 leads to a non-equilibrium increment of the tunnel current; this increment has a maximum at $V=V_0-2\Delta$ and a minimum at $V=V_0+2\Delta$. Qualitative agreement with experiment is obtained for the case of lead.

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The influence of singularities of the phonon spectrum of a superconductor on its equilibrium tunnel characteristics has been well investigated.[1] These singularities can influence also the characteristics of a non-equilibrium superconductor subjected, for example, to the action of optical radiation. Such a situation was investigated experimentally in [2], where the difference between the signals with and without radiation was measured, and it was concluded that the observed current-voltage characteristic at $V>2\Delta$ reflects the phonon structure of lead with two sharp peaks in the density of the lattice vibrations. This paper deals with the question of how the peaks in the lattice-vibration density can influence the non-equilibrium tunnel characteristics of a superconductor. It is possible that the considerations advanced below provide a qualitative explanation of the existing experimental data on lead. In addition, they can pertain to aluminum, where there is also a pronounced peak in the phonon density of states.

The gist of the effect is that in the presence of a peak in the phonon density of states at a frequency V_0, the non-equilibrium quasi-particles with energy $\epsilon > V_0+\Delta$, produced by pumping, relax more rapidly because of the possibility of emitting a phonon with energy V_0. In addition, owing to the decay of the phonon of frequency V_0 into two excitations, the influx of quasi-particles into the region $\epsilon < V_0-\Delta$ is increased. Thus, two steps appear on the non-equilibrium distribution function of the electrons, at frequencies $\epsilon = V_0 \pm \Delta$, with

$$n(\epsilon < V_0 - \Delta) > n(V_0 - \Delta < \epsilon < V_0 + \Delta) > n(\epsilon > V_0 + \Delta) .$$

Since the tunnel current depends on the electron distribution function, the presence of steps in it affects the current-voltage characteristic.

The non-equilibrium increment to the tunnel current is

$$\delta I = \frac{1}{R}\int_0^\infty \xi_\epsilon (\xi_{\epsilon+V} - \xi_{\epsilon-V}) n' d\epsilon, \qquad \xi_\epsilon = \frac{|\epsilon|}{\sqrt{\epsilon^2-\Delta^2}}\theta(\epsilon^2-\Delta^2). \tag{1}$$

The correction n' to the distribution function must be sought from the kinetic equations for the electrons and phonons. The electronic kinetic equation is analogous to the one used in [3], where it was assumed that the phonons produced in the relaxation process leave the sample. Such conditions are difficult to realize in experiment. The mean free path of a phonon of frequency $\omega \gtrsim 2\Delta$, due to scattering by electrons, is $L_1 \sim v/\Delta \sim \xi$, and for $\omega \lesssim 2\Delta$ we have analogously $\omega \lesssim 2\Delta L_2 \sim \xi \exp(\Delta/T)$ at low temperatures. If the sample thickness is $L_1 \ll d \ll L_2$, then phonons of frequency of $\omega > 2\Delta$ interact mainly with the electrons, while the $\omega < 2\Delta$ phonons interact with the thermostat. The kinetic equation for the phonons can be obtained in analogy with the electron kinetic equation. Its detailed form is given in [4]. The equation is particularly simple for the phonon distribution function N_ω at $\omega \gg \Delta$:

$$\omega N_\omega = 2 \int_\omega^\infty n_\epsilon d\epsilon . \tag{2}$$

The right-hand part contains the arrival term connected with the phonon emission by the electrons in processes in which the number of the quasiparticles is conserved. For the electron distribution function we have at $\epsilon \gg \Delta$

$$n_\epsilon \int_0^\epsilon \nu(\omega) d\omega - \int_0^{\omega_D} \nu(\omega) n_{\epsilon+\omega} d\omega - \int_\epsilon^{\omega_D} \nu(\omega) N_\omega d\omega = a \omega_D^2 \theta(\Omega - \epsilon) \tag{3}$$

where $\Omega \gg \omega_D$ is the pumping-light frequency, $\nu(\omega)$ is the density of the phonon states, λ is the electron-phonon coupling constant, $a = D(e/c)^2 A_\omega^2 (\delta/\lambda d)$ and δ is the depth of penetration of the radiation. n_ϵ in (3) stands for the distribution function averaged over the sample thickness.

With respect to $\nu(\omega)$ we make the assumption that $\nu(\omega) = 0$ at $\omega > \omega_D$ and

$$\nu(\omega) = \omega^2 + a V_0^2 f\left(\frac{\omega - V_0}{\omega_0}\right) \tag{4}$$

where the second term has the meaning of a peak in the phonon state density, smeared out over a frequency interval ω_0.

The solution for n_ϵ in the case of a smooth density of states differs from that of [3] in the presence of the large factor Ω/Δ, this being due to the "avalanche effect," since the phonons do not leave the sample and produce new quasiparticles.

In the limit $\omega_0 \gg \Delta$, when both thresholds in the electron distribution function coalesce into one at the frequency $\epsilon = V_0$, the non-equilibrium increment to the current, due to the phonon peak, takes the form

$$\frac{\delta I}{I} = -\frac{C}{V_0^7} a \omega_D^2 \Omega^2 \Delta^2 a \begin{cases} \dfrac{V - V_0}{\omega_0} \displaystyle\int_0^\infty \left| \frac{\partial f_x}{dx} \right| \frac{dx}{x} & |V - V_0| \ll \omega_0 \\ \dfrac{\omega_0}{V - V_0} \displaystyle\int_0^\infty f_x dx & |V - V_0| \gg \omega_0 \end{cases} . \tag{5}$$

The numerical constant $C \sim 1$ can be obtained by solving exactly Eqs. (2) and (3). In lead, the difference between ω_0 and Δ is not large, and the thresholds in the electron distribution function do not coalesce. As a result, a simple analysis of the dispersion of the current kernel (1) shows that the peak in the density of the phonon states at the frequency V_0 gives rise to a peak in the tunnel current at $V = V_0 - 2\Delta$ and to a dip at $V = V_0 + 2\Delta$.

The phonon density of states of lead has two peaks, at $V_1 = 4.5$ meV and $V_2 = 8.5$ meV. Current minima due to these peaks were observed in experiment [2] at $V = V_1 + 2\Delta$ and $V = V_2 + 2\Delta$, as shown in the figure. The maximum due to the second peak $V_2 - 2\Delta$ occurs in the region $V_1 + 2\Delta$ and, being the weaker, is not realized. The maximum due to the first peak $V_1 - 2\Delta$ lies in the region of frequencies lower than 2Δ and is meaningless in this situation. Matters are different when it comes to quantitative agreement with experiment. It can be shown [3] that the ratio of the amplitudes of the current dips at $V = 2\Delta$ and $V = V_1 + 2\Delta$ is

$$\frac{\delta I(2\Delta)}{\delta I(V_1 + 2\Delta)} \sim a\frac{\Delta}{T}\left(\sim\frac{V_1}{2\Delta}\right)^6 \exp\left(\frac{\Delta}{T}\right),$$

which is larger than the experimental value by two or three orders of magnitude. Of course, these results are applicable to a real superconductor only qualitatively, but the large discrepancy leaves open the question of whether the model is appropriate to the experiment.

There is one other mechanism whereby the distribution function influences the tunnel current—via changing the density of the electron states. In the phonon model of superconductivity [5] the density of states depends on the self-energy parts, which are proportional to the coupling constant and contain the distribution function. We present the result only for the relative change of the tunnel current

$$\frac{\delta I}{I} = \lambda\left(\frac{\Delta}{V_0}\right)^2 \frac{a}{\Delta^5}\, a\,\omega_D^2\,\Omega^2 \sqrt{\frac{\Delta}{T}}\,\phi(V)\exp\left(\frac{\Delta}{T}\right), \qquad (6)$$

where the function $\phi(V)$ is of the order of unity and has a dip at $V = V_0$, while the peak at $V = V_0 + 2\Delta$ and the scale of the change are of the order of the smearing of the peak in the phonon density of states.

Whereas the earlier mechanism was connected with the distribution function at high energies and was independent of temperature, this latter mechanism is governed by quasiparticles with energies the distance of which to the threshold is of the order of the temperature. Consequently, a large exponential factor appears in addition to the small coupling constant λ. However, experiment gives grounds for stating that the first, temperature-independent mechanism is more likely to operate. The indicated influence of the distribution function on the density of states may turn out to be the main mechanism, e.g., in aluminum, where the difference between Δ and V_0 is large.

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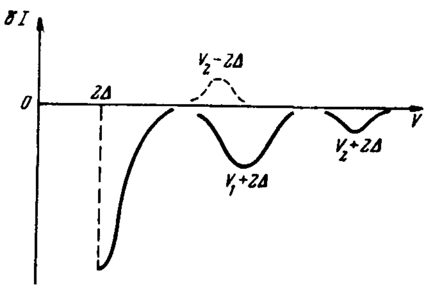

FIG. 1.

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