Creation of excess quasiparticles and nonequilibrium phase transition induced in superconductors by charged particles

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The effect of a charged-particle flux on a superconductor is investigated. The Keldysh technique is used to derive an expression for the source of excess quasiparticles in the kinetic equation. The nonequilibrium state produced by the charged particles is similar to the state resulting from the action of a broad source of optical radiation. The values of the critical flux density at which a phase transition to the normal state takes place are obtained. These values agree with the available experimental data.

INTRODUCTION

Quite a number of recent theoretical and experimental papers deal with nonequilibrium processes in superconductors acted upon by optical and laser radiation (reviewed in Refs. 1 and 2). The theoretical description is based on a system of kinetic equations in which the excess-quasiparticle source posited by Éliashberg[3] represents their creation by a monochromatic electromagnetic field. In these papers, it was shown that even under conditions of ideal heat transfer (very thin films) a phase transition to the normal state is possible, the critical superconductivity-loss powers were obtained, and the dynamics of the transition (called the nonstationary intermediate state) was investigated.[4,5]

The dynamic action of charged particles on superconductors has been studied for less, although such studies are of applied interest in connection with the use of superconductors in charged-particle accelerators, in thermonuclear reactors, and elsewhere (see, e.g., Ref. 6).

By a method of averaging over the particles, similar to that of Abrikosov and Gor'kov for paramagnetic impurities,[7] and by the Keldysh technique,[8] we obtain here a general expression for the source of the quasiparticles produced by a flux of charged particles. The shape of the source that permits solution of a number of problems is more complicated than that described by the equation for the Éliashberg source.

We obtain the critical values of the flux densities of the charged particles that return to superconductor to the normal state. The results agree with the available experimental data.[9–11]

1. FORMULATION OF PROBLEM; PRINCIPAL EQUATIONS

We consider a superconductor on whose surface is incident a flux of identical particles of density Φ, charge Ze, and velocity v. We neglect hereafter the change of the particle energy over the superconductor thickness d, assuming the latter to be much shorter than the path in the experiment.

The particles moving in the medium produce a field, which can be represented, in the nonrelativistic case of interest to us, in the form [for the gauge div $\mathbf{A}=0$ it suffices to take into account in the nonrelativistic case the scalar potential φ (Ref. 12)]

$$\mathbf{E}(\mathbf{r},t)=-\nabla\varphi(\mathbf{r},t),\quad \varphi(\mathbf{r},t)=\int\varphi(\mathbf{q},t)\exp(i\mathbf{q}\mathbf{r})\,d^3\mathbf{q},\tag{1}$$

$$e\varphi(\mathbf{q},t)=V(q)\sum_{r_a}\exp[-i\mathbf{q}(\mathbf{r}_a+\mathbf{v}t)],\tag{2}$$

$$V(q)=4\pi Ze^2/q^2\varepsilon(\mathbf{q},\mathbf{q}\mathbf{v}),\tag{3}$$

where $\mathbf{r}_a$ are the particle coordinates and $\varepsilon(\mathbf{q},\omega)$ is the dielectric constant of the medium.

The field (1) interacts with the superconductor electrons and leads to ionization and excitation. In particular, the particle field (1) creates excitations—quasiparticles—in the conduction band.

The system Hamiltonian, with allowance for the charged-particle fields, is

$$H=H_0+H_{ph}+H_{ee}+H_{imp}+V(t);\tag{4}$$

$$H_0=\sum_{\mathbf{p}}\Psi_{\mathbf{p}}^+\hat{\xi}_{\mathbf{p}}\Psi_{\mathbf{p}},$$

$$V(t)=\sum_{\mathbf{pq}}\Psi_{\mathbf{p}}^+\hat{V}(\mathbf{q},t)\Psi_{\mathbf{p}-\mathbf{q}},\tag{5}$$

where

$$\hat{\xi}_p=\begin{pmatrix}\xi_{\mathbf{p}} & \Delta\\ \Delta^* & -\xi_{\mathbf{p}}\end{pmatrix},\quad \hat{V}=\begin{pmatrix}V & 0\\ 0 & \bar{V}\end{pmatrix},\quad \begin{array}{l}V=\frac{e}{c}\mathbf{pA}+e\varphi\\ \bar{V}=\frac{e}{c}\mathbf{pA}-e\varphi\end{array}.\tag{6}$$

Here H_{ph}, H_{ee}, and H_{imp} are the Hamiltonians of the electron-phonon and Coulomb interactions and of the interaction with the impurities; ξ_p is the electron energy measured with respect to the Fermi energy; $\mathbf{A}$ and φ are the scalar and vector potentials induced in the superconductor by the charged particles (to make the approach general, we retain the potential $\mathbf{A}$); Δ is the order parameter; $\psi_{\mathbf{p}}$ and $\psi_{\mathbf{p}}^+$ are the Nambu operators in the momentum representation:

$$\Psi_{\mathbf{p}}=\begin{pmatrix}a_{\mathbf{p}\uparrow}\\ a^+_{-\mathbf{p}\downarrow}\end{pmatrix},\quad \Psi_{\mathbf{p}}^+=(a^+_{\mathbf{p}\uparrow}\,a_{-\mathbf{p}\downarrow}),$$

and $a^+_{\mathbf{p}\sigma}$ and $a_{\mathbf{p}\sigma}$ are electron creation and annihilation operators in the conduction band.

A convenient method of obtaining kinetic equations for quasiparticles is the Keldysh method[8] for Green's functions in the Gor'kov-Nambu representation (see, e.g., Refs. 13–

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15). In this case we have for the description of the system, besides the causal Green's functions G, $\widetilde{G}$, F, and F^+, also $G^{\pm}$, $\widetilde{G}^{\pm}$, etc., as defined in Refs. 8 and 13.

Interaction of the electrons with phonons, electrons, and impurities leads to the appearance of corresponding collision integrals in the kinetic equation (see, e.g., Ref. 2). We are interested in the "collision integral" $(\partial n/\partial t)_V$, corresponding to the quasiparticle source, for the quasiparticles with the charged particles [described by the fields $V(t)$ of Eq. 5].

With account taken of the expression

$$n_{\mathbf{p}} = \int_0^\infty \frac{d\omega}{2\pi i}[G^{+-}(\mathbf{p},\omega) - G^{-+}(-\mathbf{p},-\omega)] \tag{7}$$

the expression for the quasiparticle source can be written in the form

$$Q_{\mathbf{p}} = \left(\frac{\partial n}{\partial t}\right)_V = \mathrm{Im}\int_0^\infty \frac{d\omega}{2\pi}\left[\frac{\partial}{\partial t}G^{+-}(\mathbf{p},\omega) - \frac{\partial}{\partial t}G^{-+}(-\mathbf{p},-\omega)\right], \tag{8}$$

where the derivative with respect to t is understood in the sense of differentiation with respect to the "slow" time $t=(t_1+t_2)/2$, while $G(\mathbf{p},\omega) \equiv G(\mathbf{p},\omega,\mathbf{r},t)$ is taken in the mixed Wigner representation $[\mathbf{r}=(\mathbf{r}_1+\mathbf{r}_2)/2]$.

We obtain next a system of equations for the Green's functions G^{+-}, or more accurately for the matrix functions $\hat{\hat{G}}$ and $\hat{G}$

$$\hat{\hat{G}} = \begin{pmatrix} \hat{G}^{++} & \hat{G}^{-+} \\ \hat{G}^{+-} & \hat{G}^{--} \end{pmatrix}, \quad \hat{G} = \begin{pmatrix} G & iF \\ iF^+ & \widetilde{G} \end{pmatrix}. \tag{9}$$

We note first of all that the potential φ produced by the charged particles in the superconductor is formally similar to the potential of ordinary impurities[7] with density $n=\Phi/v$, the only difference being that the potentials are time-dependent [See Eq. (2)]. Actually, assume that the particles were at rest at $t=0$, and started to move at $t>0$ with velocity v. In this case, at $t=0$ the particles do not differ in any way from ordinary impurity centers randomly located in the bulk of the superconductor and having screened potentials $4\pi Ze/q^2$ $(\mathbf{q},0)$. At $t>0$ the time-dependent factor $\exp(-i\mathbf{q}\cdot\mathbf{v}t)$ and the frequency dependence of the dielectric constant $(\mathbf{q},\mathbf{q}\cdot\mathbf{v})$ are added. This can be treated not as motion of impurity centers, but as "turning on," at the instant $t=0$, a time dependence of the potential and a certain change of screening, but the "impurities" remain in their initial positions r_a.

In view of the foregoing, we can sum over the particles in the same manner as used for the impurities.[7] As a result we obtain for G the equation

$$\hat{\hat{G}}(\mathbf{p}t,\mathbf{p}'t') = \hat{\hat{G}}^0(\mathbf{p},t)\,\delta(\mathbf{p}-\mathbf{p}')\,\delta(t-t') + \sum_{\mathbf{r}_a}\iint \frac{d^3\mathbf{p}_1}{(2\pi)^3}dt_1\hat{\hat{G}}^0(\mathbf{p},t-t_1)\,\hat{\hat{V}}(\mathbf{p}-\mathbf{p}_1) \cdot \exp\{-i(\mathbf{p}-\mathbf{p}_1)(\mathbf{r}_a + \mathbf{v}t_1)\}\,\hat{\hat{G}}(\mathbf{p}_1t_1,\mathbf{p}'t'),$$
$$\hat{\hat{V}} = \begin{pmatrix} \hat{V} & 0 \\ 0 & -\hat{V} \end{pmatrix}. \tag{10}$$

2. AVERAGED EQUATIONS

In view of this analogy we shall, following Ref. 7, average Eq. (10) over the initial particle positions (at a certain instant of time $t=0$).

It is easy to verify that the ladder approximation is sufficient. Indeed, the "density" n of the charged particles, equal to Φ/v, is negligibly small, so that the parameter $p_F l$ (Ref. 7) is always very large.

The averaging yields

$$\overline{\hat{\hat{G}}(\mathbf{p}\omega,\mathbf{p}'\omega')} = \hat{\hat{G}}(\mathbf{p},\omega)\,\delta(\omega-\omega')\,\delta(\mathbf{p}-\mathbf{p}'),$$

where $\hat{\hat{G}}(\mathbf{p},\omega)$ satisfies the following Dyson equation

$$\hat{\hat{G}}(\mathbf{p},\omega) = \hat{\hat{G}}^0(\mathbf{p},\omega) + \hat{\hat{G}}^0(\mathbf{p},\omega)\,\hat{\hat{\Sigma}}(\mathbf{p},\omega), \tag{11}$$

$$\hat{\hat{\Sigma}}(\mathbf{p},\omega) = \frac{n}{(2\pi)^3}\int d^3\mathbf{q}\,\hat{\hat{V}}(q)\,\hat{\hat{G}}(\mathbf{p}-\mathbf{q},\omega-\mathbf{q}\mathbf{v}) \cdot \hat{\hat{V}}(-q)\,\hat{\hat{G}}(\mathbf{p},\omega), \tag{12}$$

$$(\omega-\xi_{\mathbf{p}})\,\hat{\hat{G}}^0(\mathbf{p},\omega) = \hat{\hat{E}}, \quad \hat{\hat{E}} = \begin{pmatrix} \hat{E} & 0 \\ 0 & -\hat{E} \end{pmatrix}. \tag{13}$$

Here $\hat{E}$ is a 2×2 unit matrix.

Using the free equations (13), we rewrite (11) in the form

$$\{\omega - \xi_{\mathbf{p}} - \hat{\hat{E}}\hat{\hat{\Sigma}}(\mathbf{p},\omega)\}\,\hat{\hat{G}}(\mathbf{p},\omega) = \hat{\hat{E}} \tag{14}$$

with a corresponding conjugate equation

$$\hat{\hat{G}}(\mathbf{p},\omega)\,\{\omega - \xi_{\mathbf{p}}^+ - \hat{\hat{\Sigma}}(\mathbf{p},\omega)\,\hat{\hat{E}}\} = \hat{\hat{E}}. \tag{15}$$

In accordance with the assumption that the density n is low, we need to calculate the function $\hat{\hat{G}}\hat{\hat{\Sigma}}\hat{\hat{E}}$ to first-order in the interaction, i.e., substitute in Eq. (14) and (15) Green's functions that satisfy the free equations but are normalized to a quasiparticle distribution function $n_{\mathbf{p}}$ as yet unknown.

The zeroth-order Green's functions can be easily obtained directly from the definitions (see, e.g., Ref. 16).

3. SOURCE OF QUASIPARTICLES PRODUCED BY A FLUX OF CHARGED PARTICLES

We now derive an expression for the quasiparticle source. Using the definition (8), we obtain for $Q_{\mathbf{p}}$ from (14) and (15)

$$Q_{\mathbf{p}} = \mathrm{Im}\int_0^\infty \frac{d\omega}{2\pi}\,[(\hat{\xi}\hat{\hat{G}} - \hat{\hat{G}}\hat{\xi}^+ + \hat{\hat{E}}\hat{\hat{\Sigma}}\hat{\hat{G}} - \hat{\hat{G}}\hat{\hat{\Sigma}}\hat{\hat{E}})^{G+-}_{\mathbf{p},\omega} - (\hat{\xi}\hat{\hat{G}} - \hat{\hat{G}}\hat{\xi}^+ + \hat{\hat{E}}\hat{\hat{\Sigma}}\hat{\hat{G}} - \hat{\hat{G}}\hat{\hat{\Sigma}}\hat{\hat{E}})^{G-+}_{-\mathbf{p},-\omega}], \tag{16}$$

where the superscripts of G^{ij} denote the corresponding block-matrix components.

To calculate Q, we subtract from the component G^{+-} of Eq. (14) the Hermitian adjoint of (15) and, taking into account the relation

$$G^{++} + G^{--} = G^{+-} + G^{-+},$$

we get

$$G^{-+}\overline{G^{+-}} - \overline{G^{-+}}G^{+-} + \overline{F^{+-}}F_+^{++} - \overline{F^{--}}F_+^{+-} - F^{+-}\overline{F_+^{++}} + F^{--}\overline{F_+^{+-}} = i[\Delta^* F^{+-} - \Delta F_+^{+-}]. \tag{17}$$

We add to this equation the expressions obtained for F_+^{+-} and F^{+-} from the corresponding components of Eq. (14) and the adjoints of (15):

$$(2\xi+\overline{G^{--}}+\overline{G^{++}})F_+^{+-}=i\Delta^*(\bar{G}^{+-}-G^{+-})+\overline{G^{+-}}F_+^{++}+\overline{F_+^{--}}G^{+-}-\overline{F_+^{+-}}G^{++}+\bar{G}^{+-}\overline{F_+^{++}}+F_+^{--}\overline{G^{+-}}-\bar{G}^{--}\overline{F_+^{+-}}, \quad (18)$$

$$(2\xi-\overline{G^{++}}-\overline{G^{--}})F^{+-}=i\Delta(\bar{G}^{+-}-G^{+-})-\overline{F^{++}}G^{+-}+\overline{F^{+-}}G^{--}-\overline{G^{+-}}F^{--}-\overline{G^{+-}}F^{++}+\overline{F^{+-}}\bar{G}^{++}-\overline{F^{--}}\bar{G}^{+-}. \quad (19)$$

In expressions (17)–(19) we have used

$$\bar{G}\equiv\bar{G}(\mathbf{p},\omega)=\frac{n}{(2\pi)^3}\int d^3\mathbf{q}V(q)V(-q)G(\mathbf{p}-\mathbf{q},\omega-\mathbf{qv}),$$

and have also taken into account that $\bar{V}\equiv-V$ in the nonrelativistic case (in the gauge div $\mathbf{A}=0$).

We obtain similar expressions also for G^{-+} and F^{-+}.

Using Eqs. (18) and (19) and the equations for F_+^{-+} and F^{-+}, we arrive at the relation

$$\mathrm{Re}\int_0^\infty d\omega[\Delta^*F^{+-}(\mathbf{p},\omega)-\Delta F_+^{+-}(\mathbf{p},\omega)]$$

$$=\mathrm{Re}\int_0^\infty d\omega[\Delta^*F^{-+}(-\mathbf{p},-\omega)-\Delta F_+^{-+}(-\mathbf{p},-\omega)].$$

Taking this relation into account, we obtain the final expression for the quasiparticle source

$$Q_\mathbf{p}=\frac{2Z^2e^4\Phi}{v}\int d^3\mathbf{q}\frac{1}{q^4|\varepsilon(\mathbf{q},\omega_\mathbf{q})|^2}\cdot\Big\{[1-n(\varepsilon_\mathbf{p})-n(\varepsilon_{\mathbf{p}-\mathbf{q}})]\delta(\varepsilon_\mathbf{p}+\varepsilon_{\mathbf{p}-\mathbf{q}}-\omega_\mathbf{q})\cdot\Big[1-\frac{\xi_\mathbf{p}\xi_{\mathbf{p}-\mathbf{q}}-|\Delta|^2}{\varepsilon_\mathbf{p}\varepsilon_{\mathbf{p}-\mathbf{q}}}\Big]+[n(\varepsilon_{\mathbf{p}-\mathbf{q}})-n(\varepsilon_\mathbf{p})]\delta(\varepsilon_\mathbf{p}-\varepsilon_{\mathbf{p}-\mathbf{q}}-\omega_\mathbf{q})\cdot\Big[1+\frac{\xi_\mathbf{p}\xi_{\mathbf{p}-\mathbf{q}}-|\Delta|^2}{\varepsilon_\mathbf{p}\varepsilon_{\mathbf{p}-\mathbf{q}}}\Big]\Big\},\quad \omega_\mathbf{q}=\mathbf{qv}. \quad (20)$$

Note that $\varepsilon(\mathbf{q},\omega)$ includes $n_\mathbf{p}$, so that (20) takes the form of the Balescu-Lenard collision integral[17] with a dielectric constant that is by itself dependent on the quasiparticle energy spectrum.

The source (20) describes the energy distribution of the quasiparticles created. Integration of $Q_\mathbf{p}$ over the momenta yields the total number of quasiparticles created

$$2\int\frac{d^3\mathbf{p}}{(2\pi)^3}Q_\mathbf{p}=\frac{4Z^2e^4\Phi}{v}\int\frac{d^3\mathbf{p}}{(2\pi)^3}\int\frac{d^3\mathbf{q}}{q^4}\frac{1}{|\varepsilon(\mathbf{q},\omega_\mathbf{q})|^2}\Big\{(1-n(\varepsilon_\mathbf{p})-n(\varepsilon_{\mathbf{p}-\mathbf{q}}))\delta(\varepsilon_\mathbf{p}+\varepsilon_{\mathbf{p}-\mathbf{q}}-\omega_\mathbf{q})\Big[1-\frac{\xi_\mathbf{p}\xi_{\mathbf{p}-\mathbf{q}}-|\Delta|^2}{\varepsilon_\mathbf{p}\varepsilon_{\mathbf{p}-\mathbf{q}}}\Big]\Big\}. \quad (21)$$

With the aid of $Q_\mathbf{p}$ we can find the total energy absorbed by the quasiparticles:

$$\frac{\partial E}{\partial t}=2\int\frac{d^3\mathbf{p}}{(2\pi)^3}\varepsilon_\mathbf{p}Q_\mathbf{p}. \quad (22)$$

At the same time, Eq. (22) should equal the energy lost by the charged particles[12]:

$$\frac{\partial E}{\partial t}=\frac{4\pi Z^2e^2\Phi}{v}\int\frac{d^3\mathbf{q}}{(2\pi)^3}\frac{\mathbf{qv}}{q^2}\frac{\mathrm{Im}[\varepsilon(q,\mathbf{qv})]}{|\varepsilon(q,\mathbf{qv})|^2}. \quad (23)$$

In fact, using the expression for the dielectric constant of the superconductor [which in our nonequilibrium situation can be easily obtained from Eq. (11) by considering the linear response to a potential of the form $V(\mathbf{r},t)=V(\mathbf{q},\omega)\times\exp(i\mathbf{qr}-i\omega t)$, just as in the calculation of the Matsubara susceptibility in Ref. 7]

$$\varepsilon(\mathbf{q},\omega)=1-\frac{8\pi e^2}{q^2}\int\frac{d^3\mathbf{p}}{(2\pi)^3}\varepsilon(\mathbf{q},\mathbf{p},\omega),$$

$$\varepsilon(\mathbf{q},\mathbf{p},\omega)_{\delta\to+0}=\frac{u_\mathbf{p}^2u_{\mathbf{p}-\mathbf{q}}^2+v_\mathbf{p}^2v_{\mathbf{p}-\mathbf{q}}^2-|\Delta|^2/2\varepsilon_\mathbf{p}\varepsilon_{\mathbf{p}-\mathbf{q}}}{\varepsilon_\mathbf{p}-\varepsilon_{\mathbf{p}-\mathbf{q}}-\omega+i\delta}(n_\mathbf{p}-n_{\mathbf{p}-\mathbf{q}})-\frac{u_\mathbf{p}^2v_{\mathbf{p}-\mathbf{q}}^2+|\Delta|^2/4\varepsilon_\mathbf{p}\varepsilon_{\mathbf{p}-\mathbf{q}}}{\varepsilon_\mathbf{p}+\varepsilon_{\mathbf{p}-\mathbf{q}}-\omega+i\delta}(1-n_\mathbf{p}-n_{\mathbf{q}-\mathbf{p}})-\frac{v_\mathbf{p}^2u_{\mathbf{p}-\mathbf{q}}^2+|\Delta|^2/4\varepsilon_\mathbf{p}\varepsilon_{\mathbf{p}-\mathbf{q}}}{\varepsilon_\mathbf{p}+\varepsilon_{\mathbf{p}-\mathbf{q}}+\omega-i\delta}(1-n_{-\mathbf{p}}-n_{\mathbf{p}-\mathbf{q}}), \quad (24)$$

we see that (23) is transformed into (22).

The expression for the source (integrated over the angles) simplifies in the most interesting case of high (compared with the Fermi velocity) velocities v of the incident particles:

$$Q(\varepsilon)=\int Q_\mathbf{p}\frac{d\Omega_\mathbf{p}}{4\pi}=\frac{dE/dt}{N(0)\omega_p^2}\sum_{i=1}^3Q_i(\varepsilon,\omega_p),$$

$$Q_1(\varepsilon,\omega)=\frac{\omega-\varepsilon+\Delta^2/\varepsilon}{[(\omega-\varepsilon)^2-\Delta^2]^{1/2}}\{1-n(\varepsilon)-n(\omega-\varepsilon)\}\theta(\omega-\varepsilon-\Delta),$$

$$Q_2(\varepsilon,\omega)=Q_3(\varepsilon,-\omega)=\frac{\varepsilon-\omega-\Delta^2/\varepsilon}{[(\varepsilon-\omega)^2-\Delta^2]^{1/2}}\{n(\varepsilon-\omega)-n(\varepsilon)\}\theta(\varepsilon-\omega-\Delta), \quad (25)$$

where $\omega_p=(4\pi n_e e^2/m)^{1/2}$ is the plasma frequency, $N(0)=mp_F/\pi^2$ is the density of states,

$$\frac{\partial E}{\partial t}=2\int Q_\mathbf{p}\varepsilon_\mathbf{p}\frac{d^3\mathbf{p}}{(2\pi)^3}=\frac{e^2Z^2\Phi\omega_p^2}{v^2}\ln\frac{2mv^2}{\omega_p}. \quad (26)$$

It was assumed in the calculation that the quasiparticle distribution function $n(\varepsilon)$ differed from zero at $\varepsilon\sim\Delta\ll\omega_p$ and the integration with respect to $|\mathbf{q}|$ was carried out by following the calculations in Ref. 18.

The source (25) is in some respects analogous to the Éliashberg source[3] that describes quasiparticle creation by a monochromatic electromagnetic field of frequency ω. Indeed, the role of the frequency ω is played in (25) by the plasma frequency ω_p. The sign in the coherent factors, however, is reversed, since the field $\mathbf{E}$ produced by the charged particles is longitudinal [see Eq. (1)].

The appearance of ω_p in (25) becomes clear if the mean value $\bar{\varepsilon}$ of the energy transferred to the quasiparticles is calculated by using the exact expression (20):

$$\bar{\varepsilon} = \int \varepsilon_p Q_p d^3\mathbf{p} \Big/ \int Q_p d^3\mathbf{p}.$$

Taking into account (26), and also

$$2\int Q_p \frac{d^3\mathbf{p}}{(2\pi)^3} = \frac{4\pi Z^2 e^2 \Phi}{v}\int \frac{d^3\mathbf{q}}{(2\pi)^3}\frac{\mathbf{qv}}{|\mathbf{qv}|}\frac{\mathrm{Im}[\varepsilon(q,\mathbf{qv})]}{q^2|\varepsilon(q,\mathbf{qv})|^2}$$
$$= \frac{e^2Z^2\Phi\omega_p}{2v^2}\ln\left(\frac{4mv^2}{\omega_p + {}^3/_5 m v_F^2}\right) \approx \frac{e^2Z^2\Phi\omega_p}{2v^2}\ln\frac{2mv^2}{\omega_p}, \qquad (27)$$

we obtain $\bar{\varepsilon} = \varepsilon_p/2$.

Since $\omega_p \gg \Delta$, the source (25) is broad in the terminology of Refs. 2 and 19.

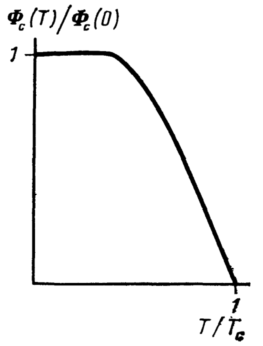

FIG. 1. Temperature dependence of the critical flux density, as given by the universal function $\beta_c(T)$ (Ref. 2).

4. CRITICAL DENSITY OF THE FLUX. COMPARISON WITH EXPERIMENT

We calculate the critical flux density Φ_c at which the order parameter vanishes. Recognizing that the source (25) is broad, we assume the kinetic equation to be homogeneous and find the connection between $n(\varepsilon)$ and Q_p from the condition that the number of quasiparticles is conserved. Integrating the kinetic equation for an isotropic and symmetric distribution function $n(\varepsilon)$ (Ref. 2) with a source (25) we obtain, just as in Ref. 20,

$$\Phi_c = \frac{n_c \hbar \omega_p \beta_c(T)}{\tau_{R\,\mathrm{eff}} R}\frac{v}{dE/dt}, \qquad (28)$$

where $n_c = mp_F\Delta_0/\pi^2\hbar^3$, $\tau_{R\,\mathrm{eff}}$ is the effective quasiparticle-recombination time, R is the multiplication coefficient, and $\beta_c(T)$ is the dimensionless critical power, a universal function of the temperatures.[2] Thus, the critical flux density is inversely proportional to the energy lost by the charged particles, to their charge, and to their velocity.

The temperature dependence of Φ_c is given by the universal function[2] $\beta_c(T)$ (see Fig. 1), so that the temperature dependences of the critical power of the optical pumping and of the charged particles are similar. Suppression of the superconductivity of lead films by an electron beam was observed in Refs. 9–11. The critical power in Ref. 9 was $2.6\cdot10^{-2}$ W. The area of the irradiated section was $S = 25\cdot10^{-6}$, so that $W_c = 10^4$ W/cm^2.

The critical flux of protons with $E = 75$ keV needed to return Nb_3Sn films to the normal state was found in Ref. 23 to be $\Phi_c = 10^{16}$ cm$^{-2}\cdot$s^{-1}. Calculation using Eq. (28) yields $\Phi_c = 1.4\cdot10^{16}$ cm$^2\cdot$s^{-1}.

Thus, the phase transition observed in Refs. 9–11 and 21 can be attributed to dynamic suppression of the superconductivity by exciting the electron system of the superconductor with a flux of incident charged particles; the nonequilibrium state produced thereby is similar to the state produced by laser pumping at a frequency ω_p.

The results also permit an estimate of the dynamic action of a neutron flux, since neutrons produce beams of high-energy ions (primary knocked-out atoms).

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