the van der Waals forces is valid. This dependence differs from the case of an infinite layer of liquid crystal where $\xi \sim H^{-2/3}$.

In conclusion, we point out that all the calculations carried out above can be generalized to the case in which the dielectric anisotropy is not small. Here, it is ture, we have considered only small departures from the equilibrium orientation. For such a calculation, the zero D functions correspond to Green's functions in a homogeneous medium. The entire change reduces to the fact that the coefficients M, ω_0 and Λ are replaced by similar expressions in which the following substitution is made everywhere in the denominators:

$$\varepsilon_3 \rightarrow \varepsilon_3 (\varepsilon_a/\varepsilon_3+1)^{1/2}. \tag{40}$$

The corresponding formulas are very involved and do not lead to any new effects. Corrections can be important only in light scattering. However, since the liquid crystals with an anisotropy of ε_a that is not small are unknown, we shall not concern ourselves with this problem here.

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[1)]Here λ_0 is the smallest characteristic length in the spectrum of the liquid crystal.

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Nonlinearity of acoustic damping in nonequilibrium superconductors

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The effect of the nonequilibrium state of an electron gas on sound absorption in a superconductor is considered. The nonequilibrium electron distribution function is assumed to be stationary and isotropic, and the phonons to be in equilibrium. In addition to the explicit dependence of the damping coefficient on the distribution function, a dependence on the correction to the order parameter exists that is associated with the nonequilibrium state. The sound intensities which create the nonequilibrium state and at which nonlinearity should be observed are estimated.

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A nonlinear dependence of the sound absorption coefficient on its intensity at temperatures below the superconducting transition has been observed in a number of experimental researches.[1] A certain decrease in the absorption above T_c has been observed in this case, which can be interpreted as an increase in the critical temperature. In addition, a deformation of the absorption curve was observed below T_c. A superconducting film in the field of a sound wave was studied in Ref. 2, and it was found that the temperature dependence of the superconducting parameters of the film in the presence of sound is supported by the theory that was developed earlier.[3] It is characteristic that the effect described in Ref. 2 was observed only for certain optimal film thicknesses, which were evidently most favorable in the sense of heat removal and the value of the sound intensity in the film.

In this connection, it should be noted that, in experiments on microwave irradiation of films, the vector of the electric field intensity is small in the film because of the smallness of the impedance at these frequencies. For this reason, microwave experiments are relatively less effective in the observation of the heating of electrons than, say, the acoustic experiments mentioned above, and experiments with laser radiation of superconducting samples.[4]

Phenomena are considered in the present work that are associated with the heating of the electron gas by the field of the sound wave. We shall call heating the isotropic change in the distribution function of the electrons which is generally not described by an effective temperature. We shall assume the phonons to be in equilibrium for the reason that near T_c a large part of the electrons remains in equilibrium and plays the role of a thermostat, together with the phonons.

The nonequilibrium situation in the case of sound propagation in a superconductor was considered by Gal'perin, Gurevich and Kozub[5] by means of an essentially anisotropic distribution function. The large magnitude of the effect in this case is ensured by the large factor kl. However, as will be shown below, at $kl < 10^3$, the heating effects appear earlier as the sound intensity increases.

1. NONLINEAR SOUND ABSORPTION COEFFICIENT

In what follows, we shall be interested in a pure superconductor $kl \gg 1$ and the case of transverse sound. In this case, in a system that is moving with the lattice, an additional exciting electromagnetic field acts on the electrons in addition to the deformation potential, in correspondence with the work of Gurevich, Lang and Pavlov.[6] The magnitude of this field must be determined from Maxwell's equations in the case of transverse sound.

Let the wave be propagated along the z axis, and let the displacement vector be equal to u_x and the vector potential A_x. Then the contribution to the Hamiltonian is

$$H' = evA_x n_x/c - iku_x pvn_x n_z, \tag{1}$$

where n is a unit vector directed along the momentum. The second term represents the deformation potential for electrons whose spectrum in the lattice differs little from the vacuum case.

The response of the electron current to the perturbation (1) is

$$j_x = \frac{3ne^2}{16mc} A_x \int_{-\infty}^{\infty} d\varepsilon \int_{-1}^{1} (1-y^2)\,dy \left[Q^R \,\mathrm{th}\frac{\varepsilon-\omega}{2T} - Q^A \,\mathrm{th}\frac{\varepsilon}{2T} + Q^a\left(\mathrm{th}\frac{\varepsilon}{2T} - \mathrm{th}\frac{\varepsilon-\omega}{2T}\right)\right] - \frac{3nev}{16} iku_x \int_{-\infty}^{\infty} d\varepsilon \int_{-1}^{1} (1-y^2)\,dy \left[q^R \,\mathrm{th}\frac{\varepsilon-\omega}{2T} - q^A \,\mathrm{th}\frac{\varepsilon}{2T} + q^a\left(\mathrm{th}\frac{\varepsilon}{2T} - \mathrm{th}\frac{\varepsilon-\omega}{2T}\right)\right]; \tag{2}$$

$$Q^R = \left[\frac{\varepsilon(\varepsilon-\omega)+\Delta^2}{\xi_\varepsilon^R \xi_{\varepsilon-\omega}^R} - 1\right] [\xi_\varepsilon^R + \xi_{\varepsilon-\omega}^R + vky + i/\tau]^{-1},$$

$$q^R = \left(\frac{\varepsilon}{\xi_\varepsilon^R} - \frac{\varepsilon-\omega}{\xi_{\varepsilon-\omega}^R}\right) \frac{y}{\xi_\varepsilon^R + \xi_{\varepsilon-\omega}^R + vky + i/\tau},$$

$$\xi_\varepsilon^R = \begin{cases} (\varepsilon^2-\Delta^2)^{1/2}\,\mathrm{sign}\,\varepsilon, & |\varepsilon| > \Delta, \\ i(\Delta^2-\varepsilon^2)^{1/2}, & |\varepsilon| < \Delta, \end{cases}$$

$\xi_\varepsilon^A = -(\xi_\varepsilon^R)$.* The expression for $Q_A(q^A)$ is obtained from $Q^R(q^R)$ by the substitution $\xi_\varepsilon^R \to \xi_\varepsilon^A$, $\xi_{\varepsilon-\omega}^R \to \xi_{\varepsilon-\omega}^A$, and the expression for $Q^a(q^a)$ by the substitution $\xi_{\varepsilon-\omega}^R \to \xi_{\varepsilon-\omega}^A$.

Near T_c, at $kl \gg 1$ and $sT/v \ll \omega \ll \Delta \ll T$ (s is the speed of sound) we have

$$j_x = \frac{3i\pi}{4}\frac{ne^2}{mc}\frac{\omega}{vk} A_x \left\{ 1 + \frac{\Delta}{2T}\ln\frac{8\Delta}{e\omega} + \frac{i\pi}{2}\frac{\Delta}{T}\frac{\Delta}{\omega} - \frac{2}{\omega}\int_{\Delta+\omega} \frac{\varepsilon(\varepsilon-\omega)+\Delta^2}{(\varepsilon^2-\Delta^2)^{1/2}((\varepsilon-\omega)^2-\Delta^2)^{1/2}} (n_\varepsilon' - n_{\varepsilon-\omega}')\,d\varepsilon - \frac{2i}{\omega}\int_{\Delta}^{\Delta+\omega} \frac{\varepsilon(\varepsilon-\omega)+\Delta^2}{(\varepsilon^2-\Delta^2)^{1/2}(\Delta^2-(\varepsilon-\omega)^2)^{1/2}} n_\varepsilon'\,d\varepsilon \right\} - eni\omega u_x, \tag{3}$$

where n_ε' is the nonequilibrium stationary contribution to the Fermi distribution function, which arises as a consequence of the nonlinearity, the source of which can be, in particular, the strong sound wave, against the background of which the damping of the weak test signal is investigated, as in the works of Fil' *et al.*[1] It is important that n_ε' is isotropic, since its anisotropic harmonics are small because of the relatively low effectiveness of the energy relaxation compared with the momentum relaxation.

The quantity responsible for the energy relaxation $\gamma \sim T^3/\omega_D^2$ is the reciprocal of the time of electron-photon interaction, and is comparable with $1/\tau$ only for very pure samples with mean free path lengths of the electrons of the order of 10^{-1} cm. Equation (3) is obtained from (2) by the formal substitution[7]

$$\mathrm{th}\,(\varepsilon/2T) \to 1 - 2n_F(\varepsilon) - 2n_\varepsilon'.$$

The absorption is determined by the imaginary part of the polarization operator, at the vertices of which are located the factors (1), and A_x is expressed in terms of u_x from Eq. (3) and the Maxwell equation $j_x = ck^2A_x/4\pi$. We shall assume the absorption Γ to be normalized to unity in the normal metal

$$\Gamma = B/(B^2+D^2); \tag{4}$$

$$B = 1 + \frac{\Delta}{2T}\ln\frac{8\Delta}{e\omega} - \frac{2}{\omega}\int_{\Delta+\omega}^{\infty} d\varepsilon \frac{\varepsilon(\varepsilon-\omega)+\Delta^2}{(\varepsilon^2-\Delta^2)^{1/2}((\varepsilon-\omega)^2-\Delta^2)^{1/2}} (n_\varepsilon' - n_{\varepsilon-\omega}'),$$

$$D = \frac{\pi}{2}\frac{\Delta}{T}\frac{\Delta}{\omega} + \frac{4}{3\pi}\frac{vk}{\omega}\left(\frac{kc}{\omega_p}\right)^2 - \frac{2}{\omega}\int_{\Delta}^{\Delta+\omega} d\varepsilon \frac{\varepsilon(\varepsilon-\omega)+\Delta^2}{(\varepsilon^2-\Delta^2)^{1/2}(\Delta^2-(\varepsilon-\omega)^2)^{1/2}} n_\varepsilon',$$

$$\omega_p^2 = 4\pi ne^2/m.$$

In the opposite limiting case $kl \ll 1$, we get in place of Eq. (3)

$$j_x = \sigma\frac{i\omega}{c}A_x - \pi\Delta\frac{\sigma}{c}[1 - 2n_F(\Delta) - 2n_\Delta']A_x - eni\omega\frac{(kl)^2}{5}u_x, \tag{5}$$

and for the absorption

$$\Gamma = \frac{1 - \Delta/2T + 2n_\Delta'}{1+F^2}, \quad F = \frac{\pi}{2}\frac{\Delta^2}{T\omega} + \frac{1}{\omega\tau}\left(\frac{kc}{\omega_p}\right)^2 - 2\pi\frac{\Delta}{\omega}n_\Delta'. \tag{6}$$

It is then seen that, thanks to the singularity at $\omega/\Delta \ll 1$, the case of the pure superconductor $kl \gg 1$ is of more interest. Here the temperature and field dependences of the absorption are stronger.

In addition to the explicit dependence of the absorption on the nonequilibrium distribution function, which is determined by Eqs. (4) and (5), there exists another source of nonlinearity, associated with the dependence of the order parameter on $n(\varepsilon)$[7,3]

$$1 = g\int_{\Delta}^{\omega_D} \frac{1-2n(\varepsilon)}{(\varepsilon^2-\Delta^2)^{1/2}}\,d\varepsilon. \tag{7}$$

Thus, the role of the sound field reduces to the creation of a nonequilibrium $n(\varepsilon)$, and with it to a change in the order parameter. Equations (4) and (6) are in fact linear absorptions against the background of given $n(\varepsilon)$ and Δ. We shall return below to the problem of the correctness of such a definition of the absorption.

For the calculation of the absorption, we need the kinetic equation connecting n'_ε with the intensity of the sound wave. The question of the kinetic equation was considered previously[7,3]; therefore, we shall not go over it in detail here; we write out the result:

$$\frac{\varepsilon}{(\varepsilon^2-\Delta^2)^{1/2}}n_\varepsilon'=\alpha_\pm\Big[\frac{\varepsilon(\varepsilon-\omega)\pm\Delta^2}{(\varepsilon^2-\Delta^2)^{1/2}((\varepsilon-\omega)^2-\Delta^2)^{1/2}}(n_{\varepsilon-\omega}-n_\varepsilon)\theta(\varepsilon-\omega-\Delta)$$
$$+\frac{\varepsilon(\varepsilon+\omega)\pm\Delta^2}{(\varepsilon^2-\Delta^2)^{1/2}((\varepsilon+\omega)^2-\Delta^2)^{1/2}}(n_{\varepsilon+\omega}-n_\varepsilon)\theta(\varepsilon-\Delta)$$
$$+\frac{\varepsilon(\varepsilon-\omega)\pm\Delta^2}{(\varepsilon^2-\Delta^2)^{1/2}((\varepsilon-\omega)^2-\Delta^2)^{1/2}}\theta(\varepsilon-\Delta)\theta(\omega-\Delta-\varepsilon)(n_{\omega-\varepsilon}+n_\varepsilon-1)\Big],$$
$$\alpha_+=\frac{8}{9\pi}\frac{(ku_xpv)^2}{\gamma vk}\frac{1}{B^2+D^2},\quad \alpha_-=\frac{2\tau(ku_xpv)^2}{15\gamma}\frac{1}{1+F^2}. \tag{8}$$

The upper sign refers to the case $kl\gg1$, the lower to the case $kl\ll1$. It is then seen that in the "dirty" limit, at $\omega\ll\Delta$, the square-root singularities are reduced by means of the coherence factor, which reduces the effect. The coherence factor in the dirty limit is the same as for the scalar potential, for the reason that in this case the fundamental contribution to the absorption is made by u_x and not A_x from (1). At $kl\gg1$, the situation is reversed.

The character of the behavior of n'_ε depends essentially on the value of α; the dependence becomes linear at $\alpha\ll(\omega/\Delta)^{1/2}$:

$$n_\varepsilon'=\frac{\alpha\omega}{4T}\Big[\frac{\varepsilon(\varepsilon-\omega)+\Delta^2}{\varepsilon((\varepsilon-\omega)^2-\Delta^2)^{1/2}}\theta(\varepsilon-\omega-\Delta)-\frac{\varepsilon(\varepsilon+\omega)+\Delta^2}{\varepsilon((\varepsilon+\omega)^2-\Delta^2)^{1/2}}\theta(\varepsilon-\Delta)\Big]. \tag{9}$$

(With an aim at treating below the case of transverse sound, we have omitted the subscript on α.) The smearing of the root singularity occurs at $(\varepsilon-\Delta-\omega)\sim\alpha^2\Delta$. At high intensities $(\omega/\Delta)^{1/2}\ll\alpha\ll(\Delta/\omega)^2$,[3] the dependence of n'_ε on α is substantially nonlinear:

$$n_\varepsilon'=\frac{\varepsilon_0}{T}f\Big(\frac{\varepsilon-\Delta}{\varepsilon_0}\Big),\quad \frac{\varepsilon_0}{\Delta}=\Big(\alpha\frac{\omega^2}{\Delta^2}\Big)^{2/5}. \tag{10}$$

There are also two modes in the sound absorption, depending on the above relations. For the values of the parameters in Eq. (4), we get, using the results obtained previously,[3] at $\alpha\ll(\omega/\Delta)^{1/2}$

$$B=1+\frac{\Delta}{2T}\ln\frac{8\Delta}{e\omega}-\alpha\frac{\Delta}{T}\Big(\frac{\Delta}{2\omega}\Big)^{1/2}\ln\frac{\omega}{\Delta\alpha^2},$$
$$D=\frac{\pi}{2}\frac{\Delta^2}{T\omega}+\frac{4}{3\pi}\frac{vk}{\omega}\Big(\frac{kc}{\omega_p}\Big)^2+\alpha\frac{\Delta}{T}\Big(\frac{\Delta}{\omega}\Big)^{1/2}a_1,\quad a_1=\frac{1}{4\sqrt{\pi}}\Gamma^2\Big(\frac14\Big), \tag{11}$$

and at $(\omega/\Delta)^{1/2}\ll\alpha\ll(\omega/\Delta)^2$, we get a strongly nonlinear mode, corresponding to (10):

$$B=1+\frac{\Delta}{2T}\ln\frac{8\Delta}{e\omega}-\frac{\Delta}{2T}\ln\frac{\varepsilon_0}{\Delta},$$
$$D=\frac{\pi}{2}\frac{\Delta^2}{T\omega}\Big(1+b_1\frac{\varepsilon_0}{\Delta}\Big)+\frac{4}{3\pi}\frac{vk}{\omega}\Big(\frac{kc}{\omega_p}\Big)^2,\quad b_1=\frac{2\pi}{(40)^{1/5}}\Big[\Gamma\Big(\frac15\Big)\sin\frac{2\pi}{5}\Big]^{-1}. \tag{12}$$

The formulas (4), (11), (12), with account of the nonequilibrium value of Δ from (7) and the expression for α from (8), in principle solve the problem of the determination of the nonlinear sound damping. We now consider the situation in which we can neglect the term proportional to $(kc/\omega_p)^2$ in the expression for D. This means that the electron and ion currents mutually compensate one another in the laboratory set of coordinates. This is possible if the sound wavelength is sufficiently large in comparison with the corresponding penetration depth of the field $(\Delta/T)^2\gg(vk/T)(kc/\omega_p)^2$. We shall assume that this inequality is valid over the entire range of real changes of Δ.

One feature of the curve of absorption as a function of temperature is the well known decrease in the absorption at $\Delta/T\gg(\omega/T)^{1/2}$ as a consequence of the Meissner screening of the eddy current which is generated by the passage of sound. Under these conditions, the nonlinear corrections to the absorption, associated with the explicit dependence of Γ on the nonequilibrium distribution function, are small. We write out the result in certain limiting cases. For this purpose, we introduce the quantity $\alpha_0=8(ku_xpv)^2/9\pi\gamma\, vk$—the value of α at $\Delta T\ll(\omega/T)^{1/2}$.

1) $(\omega/T)^{1/4}\ll\alpha_0\ll1$. Here the following temperature regions are distinguished:

a) $(\omega/T)^{1/2}\ll\Delta/T\ll1$. This is the region of rapid decrease in the absorption and the nonlinear corrections are small,

$$\Gamma=\frac{4}{\pi^2}\Big(\frac{\omega}{T}\Big)^2\Big(\frac{T}{\Delta}\Big)^4; \tag{13}$$

b) $\omega/T\alpha_0^2\ll\Delta/T\ll(\omega/T)^{1/2}$. Here we have strongly nonlinear corrections to the absorption of the order of the equilibrium correction to the value of the absorption in a normal metal,

$$\Gamma=1-\frac{\Delta}{2T}\ln\frac{8\Delta}{e\omega}-\frac{\pi^2}{4}\Big(\frac{\Delta}{T}\Big)^4\Big(\frac{T}{\omega}\Big)^2$$
$$+\frac{\Delta}{2T}\ln\Big[\alpha_0^{4/5}\Big(\frac{\Delta}{\omega}\Big)^{4/5}\Big]-\frac{b_1\pi}{2}\Big(\frac{\Delta}{T}\Big)^2\Big(\frac{\Delta}{\omega}\Big)^{6/5}\alpha_0^{2/5}; \tag{14}$$

c) $\omega/T\alpha_0^2\ll\Delta/T\ll(\omega/T)^2$. Here there is a weak nonlinear correction to the absorption,

$$\Gamma=1-\frac{\Delta}{2T}\ln\frac{8\Delta}{e\omega}-\frac{\pi^2}{4}\Big(\frac{\Delta}{T}\Big)^4\Big(\frac{T}{\omega}\Big)^2+\alpha_0\frac{\Delta}{T}\sqrt{\frac{\Delta}{2\omega}}\ln\frac{\omega}{\Delta\alpha_0^2}; \tag{15}$$

d) Upon further increase in the temperature, the increase in Δ due to stimulation of superconductivity by the sound field becomes important.[3] Δ differs from zero at $T>T_c$ all the way to $(T-T_c)/T_c\sim\omega\alpha_0/T$ and reaches the value $\Delta\sim\omega$, after which it undergoes a jump to the normal state. The absorption is also different from the value in the normal metal at $\omega\alpha_0/T\sim(T-T_c)/T_c>0$ and goes to unity by a jump upon further increase in the temperature,

$$\Gamma=1-c\omega/T\quad(c\sim1). \tag{16}$$

2) $1<\alpha_0\ll T/\omega$:

a) $(\omega/T)^{1/2}\ll\Delta/T\ll1$: the absorption is determined by Eq. (13));

b) $\omega\alpha_0^{1/2}/T\ll\Delta/T\ll(\omega/T)^{1/2}$: the absorption is determined by Eq. (14);

c) the situation is similar to 1d) at $(T-T_c)/T_c\sim\omega\alpha_0^{1/2}/T$,

$$\Delta/T\sim\omega\alpha_0^{1/2}/T,\quad \Gamma=1-c_1\frac{\omega}{T}\alpha_0^{1/2}\ln\alpha_0\qquad(c_1\sim1). \tag{17}$$

A characteristic mark of the solutions considered is the fact that the nonlinearity associated with the explicit dependence of the absorption on the nonequilibrium distribution function appears in a range of temperatures close to T_c: $\Delta/T \ll (\omega/T)^{1/2}$. The nonlinear part of the absorption is positive at first and then changes sign as one departs from T_c, because of the nonlinear increase in the Meissner screening. This can be seen, for example, from Eq. (14):

In the region of rapid decrease in the absorption $\Delta/T \gg (\omega/T)^{1/2}$, the nonlinearity that has been described is important only at $\Delta\sim T$. At high temperatures, linearity of the absorption can appear here only due to a change in Δ by the field of the wave; however, for this, a comparatively high jump level is necessary: $\alpha_0 \gg T/\omega$.

Thus, as we approach T_c from below, the absorption is linear up to $\Delta/T\sim(\omega/T)^{1/2}$; upon further increase in T, the nonlinear contribution becomes important, being first negative and then positive, in correspondence with (14) and 2b). Then the absorption again becomes less than its linear value, being located in the region $T>T_c$, in correspondence with 1d) and 2c).

The relation $\Delta(\mathbf{p})$ has been assumed to be isotropic in the calculations; in the opposite case, the results change and, in particular, the singularities of the absorption given by (14) and (15) disappear. However, the requirements on the isotropy of $\Delta(\mathbf{p})$ are not absolute. Isotropic behavior in a belt lying in a plane perpendicular to the direction of propagation of the sound is sufficient.

We have determined the absorption as the imaginary part of the polarization operator; we shall show that in such a method of determination, the absorption is proportional to the energy radiated by the electron system in the form of thermal phonons. The kinetic equation is written with the help of the Green's function integrated over ξ[7]:

$$I(\varepsilon)=\int\frac{e}{c}vA\,(g_{\varepsilon-\omega\varepsilon}-g_{\varepsilon\varepsilon+\omega}+g_{\varepsilon+\omega\varepsilon}-g_{\varepsilon\varepsilon-\omega})\frac{dO_p}{4\pi}. \tag{18}$$

The inelastic collision integral is on the left side and has a Bloch form at $\Delta=0$. We introduce the function

$$\psi(\varepsilon)=\int_{-\infty}^{\varepsilon}\frac{d\varepsilon_1}{4\pi i}\int\frac{e}{c}vAg_{\varepsilon_1\varepsilon_1-\omega}\frac{dO_p}{4\pi}, \tag{19}$$

Then the absorption $\Gamma\sim\operatorname{Im}\psi(\infty)$. We integrate (14) with account of the symmetry of the functions g:

$$\int_{-\infty}^{\varepsilon_1}I(\varepsilon_2)\frac{d\varepsilon_2}{4\pi i}=2\psi(\varepsilon_1+\omega)-2\psi(\varepsilon_1).$$

Integrating the resultant expression over $d\varepsilon_1$ in the limits from $-\infty$ to ε and extending ε to $+\infty$, we get, after simple transformations,

$$\psi(\infty)=-\frac{1}{2\omega}\int_{-\infty}^{\infty}\frac{d\varepsilon}{4\pi i}\varepsilon I(\varepsilon).$$

The collision integral can be tentatively represented in the form

$$I(\varepsilon)=\int_{-\infty}^{\infty}d\varepsilon_1F[N(\varepsilon_1-\varepsilon),n(\varepsilon),n(\varepsilon_1)], \tag{20}$$

where $N(\varepsilon_1-\varepsilon)$ is the phonon distribution function. The integrand of (20) does not change as a result of the substitution $\varepsilon\to-\varepsilon$, $\varepsilon_1\to-\varepsilon_1$. Then, with account of this, we have

$$\Gamma\sim\operatorname{Im}\psi(\infty)=\frac{1}{2\omega}\int_{-\infty}^{\infty}\frac{d\varepsilon}{4\pi}\int_{-\infty}^{\infty}(\varepsilon-\varepsilon_1)F[N(\varepsilon_1-\varepsilon),n(\varepsilon),n(\varepsilon_1)]d\varepsilon_1.$$

The quantity at the right is the energy radiated in the form of thermal phonons, which proves the assumption made above.

It is interesting to compare the nonlinear parameter α_0 with the corresponding parameter of the work of Gal'perin *et al.*,[5] which corresponds to another nonlinearity mechanism: $(ku)^2(kl)^4\sim(ku)^2\varepsilon_F^2/\gamma vk$, $kl\sim(\varepsilon_F\tau/\gamma)^{1/5}$. The high power of the root makes the estimate practically insensitive to any real change in τ. We then obtain the result that the heating parameter contains a nonlinearity at $kl<10^3$.

We estimate the required intensity of the sound wave at which $\alpha_0\sim1$. We introduce the energy flux density $W\sim\rho s\omega^2u_x^2$, where ρ is the density of the metal. Then, for $\omega\sim10^9$–10^{10} sec^{-1}, $W\sim10^{-2}$ W/cm^2. Heating effects are observed experimentally at $|T_c-T|<10^{-2}$ deg, which under our conditions corresponds to the estimates $\Delta/T<(\omega/T)^{1/2}$, $\omega T>s/v$. The Meissner screening increases at lower frequencies, which narrows the temperature interval of the existence of the effect.

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