Fluctuation conductivity of superconductors of the second kind above T_c in the presence of strong electric and magnetic fields

O. D. Cheĭshvili

Institute of Physics, Georgian Academy of Sciences
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The fluctuation correction to the conductivity above the superconducting transition temperature is investigated for the case of dirty superconductors in the presence of constant electric and magnetic fields. The fluctuation conductivity is calculated for a bulk sample in mutually parallel and mutually perpendicular electric and magnetic fields. The fluctuation conductivity of a thin film in mutually perpendicular fields is also calculated. The asymptotic values of the conductivities in strong fields are obtained.

It is well known that the existence of fluctuating electron pairs above the superconducting transition temperature leads to the result that a normal metal acquires near the transition temperature an additional conductivity, which is strongly temperature dependent (the so-called paraconductivity). The first microscopic calculation of the paraconductivity was made in the article by Aslamazov and Larkin.[1] This same question was investigated on the basis of the time-dependent Ginzburg-Landau equation in the articles by Abrahams and Woo,[2] and by Schmid.[3] These calculations were found to be in good agreement with the experimental data in the articles by Glover,[4] and by Strongin et al.[5]

The phenomenon of paraconductivity in the absence of a magnetic field was investigated in[1-3]. The electric field was assumed to be sufficiently weak so that it would not lead to depairing of the fluctuating electron pairs. The criterion for the weakness of the electric field will be discussed below. In the case of ordinary (not paramagnetic) alloys, Maki[6] and Thompson[7] showed that it is necessary to correct the theoretical results[1-3] by taking the so-called "nonregular" terms[8] into consideration. The nonlinear dependence of the paraconductivity on the electric field has been studied by a number of authors.[9-14] The fluctuation conductivity in the presence of a constant magnetic field has also been investigated.[14-17]

The fluctuation conductivity of dirty superconductors above T_c in strong electric and magnetic fields is investigated in the present article, taking into account only the fluctuations of the Aslamazov and Larkin type.[1)] We shall calculate the electric current according to the formula

$$\mathbf{J}(\mathbf{r},t)=\frac{e}{m}\lim_{\mathbf{r}\to\mathbf{r}',t\to t'}(\partial_{\mathbf{r}}-\partial_{\mathbf{r}'}^{*})G^{+}(\mathbf{r},t;\mathbf{r}',t'), \tag{1}$$

where $\partial_{\mathbf{r}} = \partial/\partial\mathbf{r} - ie\mathbf{A}$, $\mathbf{A} = (0, xH, 0)$ is the vector potential of the constant magnetic field, and $G^{+}(x, x')$ is the Green's function of the electron in constant electric and magnetic fields. According to Keldysh,[19]

$$G^{+}(x,x')=i\langle\psi^{+}(x')\psi(x)\rangle.$$

To first-order in the fluctuations, the electron Green's function corresponds to a graph (see Fig. 1). The solid lines in Fig. 1 represent the Green's functions of the electron in the normal state. The heavy, wavy line represents the Cooper vertex function in constant electric and magnetic fields.

In what follows it will be necessary for us to expand the electron Green's functions and the vertex function in this graph in powers of the electromagnetic field $\mathbf{A}_0(t) = -\mathbf{E}t$. We shall take the constant magnetic field into account exactly. It is necessary to emphasize that the Aslamazov and Larkin contributions to the conductivity give graphs involving the interaction of the electromagnetic field with fluctuating pairs of electrons. However, the graphs containing the interaction of the electromagnetic field with the normal electrons give the Maki and Thompson contributions to the conductivity. As noted earlier, we shall confine our attention to graphs of the Aslamazov and Larkin type; therefore, the electric field is included only in the Cooper vertex function.

Using Keldysh's technique,[19,20] the fluctuation correction to the current, corresponding to the graph (Fig. 1), can be written in the following form:

$$\Delta\mathbf{J}_{\omega_0}=i\frac{e}{4m}\lim_{\mathbf{r}\to\mathbf{r}'}(\partial_{\mathbf{r}}-\partial_{\mathbf{r}'}^{*})\int d\mathbf{r}_1\,d\mathbf{r}_2\frac{d\omega}{2\pi}\frac{d\varepsilon}{2\pi}\,\mathrm{th}\frac{\varepsilon}{2T}$$
$$\times[\langle G_{\varepsilon}^{R}(\mathbf{r},\mathbf{r}_1)G_{\omega-\varepsilon}^{A}(\mathbf{r}_2,\mathbf{r}_1)G_{\varepsilon-\omega_0}^{R}(\mathbf{r}_2,\mathbf{r}')\rangle \tag{2}$$
$$-\langle G^{A}(\mathbf{r},\mathbf{r}_1)G_{\omega-\varepsilon}^{R}(\mathbf{r}_2,\mathbf{r}_1)G_{\varepsilon-\omega_0}^{A}(\mathbf{r}_2,\mathbf{r}')\rangle]K_{\omega,\omega-\omega_0}(\mathbf{r}_1,\mathbf{r}_2).$$

Here the angle brackets indicate averaging over impurities,[21] G^R and G^A denote the retarded and advanced Green's functions, respectively, in the presence of a constant magnetic field, and $K_{\omega,\,\omega-\omega_0}(\mathbf{r}_1,\mathbf{r}_2)$ is the so-called thermodynamic Cooper vertex function (see[21]).

The expansion of the Cooper vertex function in powers of the variable vector potential is shown graphically in Fig. 2. The thin wavy lines on these diagrams represent the Cooper vertex functions in a constant magnetic field, and the dotted lines correspond to the interaction $e\mathbf{p}\cdot\mathbf{A}_0/m$ of an electron with the variable electromagnetic field. Averaging over impurities[20] is also assumed. Only odd powers in $e\mathbf{p}\cdot\mathbf{A}_0/m$ are present in the diagrams. It is not difficult to see that the even powers in expression (2) for the current do not give any contribu-

FIG. 1

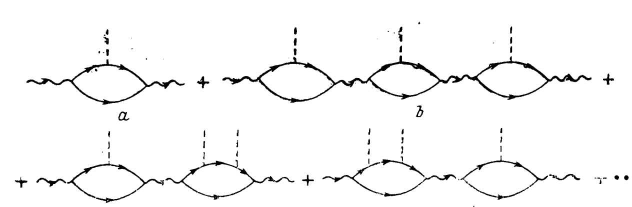

FIG. 2

tion. They drop out because of the integration over the momenta.

By analytic continuation of the vertex function, we see that the result is analogous in structure to the expression for the Green's functions of an electron in a variable electromagnetic field (see[8]). As an example let us write down the explicit form of graph c in Fig. 2, pertaining to the third-order correction in $\mathbf{A}_0$:

$$\begin{aligned} K_{\omega,\omega-\omega_0}(\mathbf{r},\mathbf{r}') &= \Pi^{(1)}(\mathbf{r}_1,\mathbf{r}_2,\omega_1)\,\Pi^{(2)}(\mathbf{r}_3,\mathbf{r}_4,\omega_2,\omega_3) \\ &\times\Big[\operatorname{cth}\frac{\omega}{2T}K_\omega^R(\mathbf{r},\mathbf{r}_1)K^R_{\omega-\omega_1}(\mathbf{r}_2,\mathbf{r}_3)K^R_{\omega-\omega_1-\omega_2-\omega_3}(\mathbf{r}_4,\mathbf{r}') \\ &-\operatorname{cth}\frac{\omega-\omega_1-\omega_2-\omega_3}{2T}K_\omega^A(\mathbf{r},\mathbf{r}_1)K^A_{\omega-\omega_1}(\mathbf{r}_2,\mathbf{r}_3)K_{\omega-\omega_1-\omega_2-\omega_3}(\mathbf{r}_4,\mathbf{r}') \\ &-\left(\operatorname{cth}\frac{\omega}{2T}-\operatorname{cth}\frac{\omega-\omega_1}{2T}\right)K_\omega^A(\mathbf{r},\mathbf{r}_1)K^R_{\omega-\omega_1}(\mathbf{r}_2,\mathbf{r}_3)K^R_{\omega-\omega_1-\omega_2-\omega_3}(\mathbf{r}_4,\mathbf{r}') \\ &-\left(\operatorname{cth}\frac{\omega-\omega_1}{2T}-\operatorname{cth}\frac{\omega-\omega_1-\omega_2-\omega_3}{2T}\right)K_\omega^A(\mathbf{r},\mathbf{r}_1)K^A_{\omega-\omega_1}(\mathbf{r}_2,\mathbf{r}_3)K^R_{\omega-\omega_1-\omega_2-\omega_3}(\mathbf{r}_4,\mathbf{r}')\Big\}. \end{aligned} \quad (3)$$

Here the integration with respect to internal coordinates and with respect to the frequencies ω_i is understood to be subject to the condition $\omega_1 + \omega_2 + \omega_3 = \omega_0$; $\Pi^{(1)}$ and $\Pi^{(2)}$ correspond to loop diagrams with one and two electromagnetic lines; more precisely, $\Pi^{(1)}$ is the sum of loops a and b, and $\Pi^{(2)}$ is the sum of loops c, d, and e (Fig. 3).

As estimates indicate, loops with more than two electromagnetic lines are not essential in view of the smallness of the parameter $l(T - T_c)v_0 \ll 1$ (l denotes the mean free path).

The expression for the Cooper vertex function $K_{\omega,\,\omega-\omega_0}$ to arbitrary order in the electromagnetic field is written down in analogy to expression (3) and does not require any additional explanation. In formula (3) K^R and K^A denote the retarded and advanced vertex functions in a constant magnetic field. Maki showed[22] that, for very dirty superconductors the vertex function in a constant magnetic field is diagonalized in the system of eigen wavefunctions of the following equation:

$$-\tilde{\partial}_r^2\varphi_{nkq}(\mathbf{r}) = \varepsilon_{nq}\varphi_{nkq}(\mathbf{r}), \quad (4)$$

where $\tilde{\partial}_\mathbf{r} = \partial/\partial\mathbf{r} - 2ie\mathbf{A}$, $\mathbf{A} = (0, xH, 0)$, $\epsilon_{nq} = 2eH(2n + 1) + q^2$ are the eigenvalues, and the wave functions have the form

$$\varphi_{nkq}(\mathbf{r}) = \frac{1}{2\pi(2^n n!\pi^{1/2}\lambda)^{1/2}}\exp\left\{iky+iqz-\frac{1}{2}\left(\frac{x-x_0}{\lambda}\right)^2\right\}H_n\left(\frac{x-x_0}{\lambda}\right), \quad (5)$$

where $x_0 = \lambda^2 k$, $\lambda = (2eH)^{-1/2}$, and the $H_n(x)$ are the Hermite polynomials.

Thus, in the case of dirty superconductors one can write down the following expression for the Cooper vertex function in the representation of Landau quantum numbers:

$$\begin{aligned} K_\omega^R(n,q) &= \nu^{-1}\left[\ln\frac{T}{T_c(0)}+\psi\left(\frac{1}{2}+\frac{D\varepsilon_{nq}-i\omega}{4\pi T}\right)-\psi\left(\frac{1}{2}\right)\right]^{-1}, \\ K_\omega^A(n.\ q) &= [K_\omega^R(n,\ q)]^*. \end{aligned} \quad (6)$$

Here $\nu = mp_0/2\pi^2$ is the electron density of states, $D = v_0 l/3$ is the coefficient of diffusion, $T_c(0)$ is the transition temperature when no magnetic field is present, and $\psi(x)$ is the logarithmic derivative of the $\Gamma(x)$ function.

Making the calculations, similar to Usadel's work,[15] we obtain an expression for the loop with one electro-

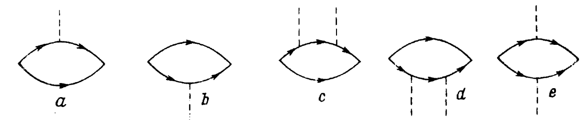

FIG. 3

magnetic line in the presence of a constant magnetic field in the form

$$\begin{aligned} \int d\mathbf{r}\,d\mathbf{r}'\,\Pi^{(1)}(\mathbf{r},\mathbf{r}',\omega)\varphi^*_{nkq}(\mathbf{r})\varphi_{n'k'q'}(\mathbf{r}) &= \Pi^{(1)}_{nn'}(q,\omega)\delta(k-k')\delta(q-q') \\ &= -\frac{e\nu D}{2\pi T}L_{nn'}(q)\delta'(\omega)\int d\mathbf{r}(\tilde{\partial}_\mathbf{r}^*-\tilde{\partial}_{\mathbf{r}'})\mathbf{E}\varphi^*_{nkq}(\mathbf{r})\varphi_{n'k'q'}(\mathbf{r}')|_{\mathbf{r}=\mathbf{r}'}, \\ L_{nn'}(q) &= \frac{\psi(1/2+D\varepsilon_{nq}/4\pi T)-\psi(1/2+D\varepsilon_{n'q}/4\pi T)}{D(\varepsilon_{nq}-\varepsilon_{n'q})/4\pi T} \end{aligned} \quad (7)$$

The following expression corresponds to the loop with two electromagnetic lines:

$$\begin{aligned} \int d\mathbf{r}\,d\mathbf{r}'\Pi^{(2)}(\mathbf{r},\mathbf{r}',\omega,\omega')\varphi^*_{nkq}(\mathbf{r})\varphi_{n'k'q'}(\mathbf{r}') &= \Pi^{(2)}_n(q,\omega,\omega')\delta(k-k')\delta(q-q')\delta_{nn'} \\ &= (eE)^2\frac{\nu D}{\pi T}\psi^{(1)}\left(\frac{1}{2}+\frac{D\varepsilon_{nq}}{4\pi T}\right)\delta'(\omega)\delta'(\omega')\delta(k-k')\delta(q-q')\delta_{nn'}, \end{aligned} \quad (8)$$

where $\psi^{(1)}(x)$ is the derivative of the $\psi(x)$ function.

After similar calculations expression (2) for the current can be rewritten in the following form:

$$\begin{aligned} \Delta\mathbf{J}_{\omega_0} &= \frac{e\nu D}{4\pi T}(\tilde{\partial}_\mathbf{r}-\tilde{\partial}_{\mathbf{r}'}^*)\sum_{nn'}\int dk\,dq\frac{d\omega}{2\pi}\varphi_{nkq}(\mathbf{r})\varphi^*_{n'kq}(\mathbf{r}')|_{\mathbf{r}=\mathbf{r}'} \\ &\times L_{nn'}(q)K_{\omega,\omega-\omega_0}(nq,n'q), \end{aligned} \quad (9)$$

where $K_{\omega,\,\omega-\omega_0}(nq, n'q)$ is the sum of all the diagrams for the vertex function (Fig. 2). In this formula nq and n'q are the Landau quantum numbers of the external Cooper functions in the diagrams. The fact that the vertex function is given by

$$K_{\omega,\omega-\omega_0}(nkq,\ n'k'q') = K_{\omega,\omega-\omega_0}(nq,\ n'q)\delta(k-k')\delta(q-q')$$

has been taken into consideration in deriving this formula. As one can easily see, this assertion follows from formulas (6), (7), and (8). From formula (9) one can easily see that the current doesn't depend on the spatial coordinates. Returning to the time representation, we have the following expression for the current:

$$\begin{aligned} \Delta\mathbf{J}(t) &= \frac{e\nu D}{4\pi T}(\tilde{\partial}_\mathbf{r}-\tilde{\partial}_{\mathbf{r}'}^*)\sum_{nn'}\int dk\,dq\,d\omega_0\frac{d\omega}{2\pi}\varphi_{nkq}(\mathbf{r})\varphi^*_{n'kq}(\mathbf{r}')|_{\mathbf{r}=\mathbf{r}'} \\ &\times L_{nn'}(q)K_{\omega,\omega-\omega_0}(nq,n'q)e^{-i\omega_0 t}. \end{aligned} \quad (10)$$

Expression (10) is a function of the time. However, it should be noted that the time-dependent part of expression (10) is cancelled by the current arising from the "regular"[8] part of the Maki and Thompson graphs. Therefore, in what follows we shall everywhere present the time-independent part of the current, $\Delta\mathbf{J}(0) \equiv \Delta\mathbf{J}$. A similar situation was noted by Aslamazov and Larkin[1] in connection with a calculation of the fluctuation current to the linear approximation in the electric field.

After introducing the general relationships, let us go on to a calculation of the fluctuation correction to the conductivity of a bulk sample in the presence of mutually parallel and mutually perpendicular electric and magnetic fields, and the conductivity of a thin layer when the magnetic field is applied perpendicular to the plane of the layer.

Let us consider the conductivity of a bulk sample in the case when $\mathbf{E} \parallel \mathbf{H} \parallel \mathbf{z}$. Substituting the wave functions

(5) into expression (10) for the current, it is not difficult to find that

$$\Delta J_{\parallel}=i\frac{evD}{2\pi T(2\pi\lambda)^2}\sum_n\int dq\,d\omega_0\frac{d\omega}{2\pi}qL_{nn}(q)K_{\omega,\omega-\omega_0}(nq,nq). \quad (11)$$

It is clear from formulas (5) and (7) that

$$\Pi^{(1)}_{nn'}(q,\omega)=i\frac{evD}{\pi T}Eq\psi^{(1)}\left(\frac{1}{2}+\frac{D\varepsilon_{nq}}{4\pi T}\right)\delta'(\omega)\delta_{nn'}. \quad (12)$$

Thus, in this case the interaction of the fluctuating pairs with the variable electromagnetic field is determined (in the Landau representation) by expressions (8) and (12). Using this result, and also formulas (6) and (11), and after carrying out the integration to arbitrary order in the electric field, we obtain the following result for the fluctuation conductivity:

$$\Delta\sigma_{\parallel}^{(3)}=\frac{e^2HTD^{1/2}}{2\pi}\sum_{n,m=0}^{\infty}(-)^m\frac{(6m+1)!!}{2^{5m}3^m m!}(eE)^{2m}D^m\left(\frac{L_n}{4\pi T\Lambda_n}\right)^{3m+1/2}$$

$$=\frac{e^2H}{4\pi^3}\left(\frac{D}{T}\right)^{1/2}\sum_n\int_0^{\infty}x^2\exp\left\{-\frac{\Lambda_n}{L_n}x^2-\frac{(eE)^2D}{12(4\pi T)^3}x^6\right\}dx; \quad (13)$$

$$\Lambda_n=\ln(T/T_c(0))+\psi(1/2+(2n+1)eHD/2\pi T)-\psi(1/2),$$
$$L_n=\psi^{(1)}(1/2+(2n+1)eHD/2\pi T).$$

It is not difficult to obtain the limit $H\to 0$ from formula (13). Substituting $L_n\approx\pi^2/2$ and $\Lambda_n\approx\ln(T/T_c(0))+\pi(2n+1)eHD/4T$ into (13) and summing over the principal quantum number n, we obtain

$$\Delta\sigma_{\parallel}^{(3)}=\frac{e^3H}{16}\left(\frac{D}{2T}\right)^{1/2}\int_0^{\infty}x^2\exp\left\{-x^2\ln\frac{T}{T_c(0)}-x^6\frac{\pi^3(eE)^2D}{12(8T)^3}\right\}\mathrm{sh}^{-1}\frac{\pi eHDx^2}{4T}dx. \quad (14)$$

Expression (14) coincides with the result of [23], which is valid in the Ginzburg-Landau regime, $eHD\ll\pi T$. For $H=0$ we obtain from Eq. (14) the well known result [10] for the nonlinear (in the electric field) conductivity in zero magnetic field:

$$\Delta\sigma_{\parallel}^{(3)}(H=0)=\frac{2}{\sqrt{\pi}}\Delta\sigma_0^{(3)}\int_0^{\infty}\exp\left\{-x^2-x^6\frac{E^2}{E_c^2}\right\}dx, \quad (15)$$

$$\Delta\sigma_0^{(3)}=\frac{e^2T_c(0)}{8(2\pi D(T-T_c(0)))^{1/2}}\qquad E=\frac{32\sqrt{6}(T-T_c(0))^{3/2}}{\pi^{3/2}eD^{1/2}}=E\left(\frac{T-T_c(0)}{T_c(0)}\right)^{3/2} \quad (16)$$

As one can easily see from formula (13), the limiting value as $E\to 0$ in an arbitrary magnetic field coincides with the result of Usadel's work. [15]

Let us consider the case of strong electric and magnetic fields, $eHD\gg\pi T$, and let $T\gtrsim T_c(H)$ ($T_c(H)$ is the transition temperature in the presence of a magnetic field). Then one can replace the summation in formula (13) by the first term $n=0$. For temperatures close to $T_c(H)$ we have

$$\frac{\Lambda_0}{L_0}\approx\frac{1-\rho(H)\psi^{(1)}(1/2+\rho(H))}{\psi^{(1)}(1/2+\rho(H))}\frac{T-T_c(H)}{T_c(H)}\equiv\gamma(H), \quad (17)$$
$$\rho(H)=eHD/2\pi T_c(H).$$

Substituting (17) into (13), we obtain after simple transformations

$$\Delta\sigma_{\parallel}^{(3)}=\frac{e^3H}{4\pi^3}\left(\frac{D}{T_c(H)}\right)^{1/2}[\gamma(H)]^{-3/2}\int_0^{\infty}x^2\exp\left\{-x^2-\frac{E^2}{E_{c1}^2}x^6\right\}dx, \quad (18)$$

$$E_{c1}=\frac{16\sqrt{3}}{eD^{1/2}}[\pi T_c(H)\gamma(H)]^{3/2}. \quad (19)$$

In the case of strong electric fields, $E_{c1}\ll E$, we find

$$\Delta\sigma_{\parallel}^{(3)}(H)\approx 2e^2HT_c(H)/\sqrt{3}\pi E. \quad (20)$$

However, in the case $H\gtrsim H_{c2}(T)$ and in the presence of a strong electric field,

$$E\gg E_{c2}=4\sqrt{6}e^{1/2}D[H-H_{c2}(T)]^{3/2} \quad (21)$$

we obtain

$$\Delta\sigma_{\parallel}^{(3)}(T)\approx 2e^2TH_{c2}(T)/\sqrt{3}\pi E. \quad (22)$$

It is clear from formulas (20) and (22) that the electric field dependence of the fluctuation conductivity for a three-dimensional metal in the presence of strong electric and magnetic fields is the same as for the one-dimensional case in a strong electric field. The condition imposed on the electric field is less stringent than in the case without a magnetic field. In fact, it is clear from formulas (16) and (19) that in order of magnitude

$$E_{c1}\sim E_{c0}\left(\frac{T-T_c(H)}{T_c(0)}\right)^{3/2}\left(\frac{\pi T_c(H)}{eHD}\right)^{1/2} \quad (23)$$

for $eHD/\pi T_c(H)\gg 1$.

Let us consider the case of mutually perpendicular electric and magnetic fields. Let $\mathbf{E}\parallel x$ and $\mathbf{H}\parallel z$; then it is clear from Eqs. (5), (7), and (10) that

$$\Delta J_{\perp}=\frac{evD}{2T(2\pi\lambda)^3}\sum_{nn'}[(2n)^{1/2}\delta_{n-1,n'}-(2(n+1))^{1/2}\delta_{n,n'-1}]$$
$$\times\int dq\,d\omega_0\frac{d\omega}{2\pi}L_{nn'}(q)K_{\omega,\omega-\omega_0}(nq,n'q), \quad (24)$$

$$\Pi^{(1)}_{nn'}(q,\omega)=-\frac{evDE}{2\pi T\lambda}L_{nn'}(q)\delta'(\omega)[(2n)^{1/2}\delta_{n-1,n'}-(2(n+1))^{1/2}\delta_{n,n'-1}]. \quad (25)$$

For mutually perpendicular electric and magnetic fields we shall, from the very beginning, confine our attention to the case of strong magnetic fields, $eHD/2\pi T\gg 1$. Then one can keep the lowest Landau levels in the summation in Eq. (24). Using formulas (8), (24), and (25), after the summation of all vertex diagrams (Fig. 2) we obtain the following result for the fluctuation conductivity of a three-dimensional sample:

$$\Delta\sigma_{\perp}^{(3)}=\frac{2}{\sqrt{\pi}}\Delta\sigma_{0\perp}^{(3)}\int_0^{\infty}\exp\left\{-x^2-\frac{E^2}{E_{c\perp}^2}x^6\right\}dx. \quad (26)$$

However, in the case of a thin film of thickness d, when the magnetic field is perpendicular to the plane of the film, the fluctuation conductivity has the form

$$\Delta\sigma_{\perp}^{(2)}=\Delta\sigma_{0\perp}^{(2)}\int_0^{\infty}\exp\left\{-x-\frac{E^2}{E_{c\perp}^2}x^3\right\}dx. \quad (27)$$

In these formulas

$$\Delta\sigma_{0\perp}^{(3)}=\frac{e^2\lambda T}{3\pi D\Lambda_0^{1/2}},\qquad\Delta\sigma_{0\perp}^{(2)}=\frac{2e^2\lambda^2T}{3\pi dD\Lambda_0},$$
$$E_{c\perp}=\left(\frac{6}{2-\ln 3}\right)^{1/2}\frac{D\Lambda_0^{3/2}}{e\lambda^3},$$
$$\Lambda_0=\ln(T/T_c(0))+\psi(1/2+eHD/2\pi T)-\psi(1/2).$$

It is easy to verify that for $E=0$ formulas (26) and (27) go over into the corresponding formulas of Usadel's work. [15]

In the case of a strong electric field, $E\gg E_{c\perp}$, and for temperatures $T\gtrsim T_c(H)$ we find

$$\Delta\sigma_{\perp}^{(3)}\approx\frac{\Gamma(1/6)}{9\pi^{3/2}}\left(\frac{6}{2-\ln 3}\right)^{1/6}\frac{e^2T_c(H)}{D^{1/3}(eE)^{1/3}}, \quad (28)$$

$$\Delta\sigma_{\perp}^{(2)} \approx \frac{2\Gamma(^1/_3)}{9\pi}\left(\frac{6}{2-\ln 3}\right)^{1/3} \frac{e^2 T_c(H)}{dD^{1/3}(eE)^{2/3}}. \quad (29)$$

It should be noted that in these asymptotic expressions the electric field dependence of the conductivity is the same as in the case when no magnetic field is present. The magnetic field dependence manifests itself only through $T_c(H)$, and the conductivity—in the two-dimensional case as well as in the three-dimensional case—decreases with decreasing transition temperature and increasing magnetic field.

[1)]It should be noted that the fluctuations of the Maki and Thompson type are strongly suppressed in the presence of magnetic fields [[18]].

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