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MAGNETIC FIELD ENHANCEMENT OF THE c-AXIS RESISTIVITY PEAK NEAR T_c IN LAYERED SUPERCONDUCTORS

V.V.Dorin[†], *R.A.Klemm*[†], *A.A.Varlamov*[†], *A.I.Buzdin*[*†], *D.V.Livanov*[+]

[†]*Materials Science Division, Argonne National Laboratory*
9700 S. Cass Ave., Argonne, IL 60439

[*]*Institute for High Pressure Physics RAS*
142092 Troitsk, Russia

[+]*Moscow Institute of Steel and Alloys,*
117936 Moscow, Russia

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The contributions to the c-axis conductivity from fluctuations of the normal quasiparticle density of states are opposite in sign to the Aslamazov–Larkin and Maki–Thompson contributions, leading to a peak in the overall c-axis resistivity $\rho_c(T)$ above T_c. In a magnetic field $\mathbf{H}||\hat{c}$, this peak increases in magnitude and is shifted to lower T by an amount proportional to H^2 for weak fields and to H for strong fields. Our results are discussed in regard to recent experiments with $YBa_2Cu_3O_{7-\delta}$ and $Bi_2Sr_2CaCu_2O_{8+\delta}$.

One of the main characteristics of a superconductor is the temperature T dependence of the resistivity in the vicinity of the superconducting transition temperature T_c. Recently, the resistivities ρ of many high-T_c cuprates have been studied by numerous groups [1-7]. Reproducible results on untwinned samples of $YBa2Cu_3O_{7-\delta}$ (YBCO) [2-4] and on $Bi_2Sr_2CaCu_2O_{8+\delta}$ (BSCCO) [1,5] have been obtained. In BSCCO, $\rho_c(T)$ exhibits a peak, which increases in magnitude with decreasing oxygen concentration [6]. In addition, a rather weak magnetic field $\mathbf{H}||\hat{c}$ causes the relative magnitude of the peak to increase, and its position to shift dramatically to lower $T > T_c(H)$ values [5]. On the other hand, fully-oxygenated YBCO appears metallic along all three crystal axis directions, although oxygen-deficient YBCO can exhibit a peak in $\rho_c(T)$ [7].

Recently, it was proposed [8] that such a peak in the zero-field $\rho_c(T)$ could arise from superconducting fluctuations, and the calculations [8] were found to be in agreement with experiments on epitaxially-grown thin films of BSCCO [6]. The main idea of this explanation is that the Aslamazov–Larkin (AL) fluctuation conductivity contribution in the c-axis direction is *weak* in the '2D' regime above the dimensional crossover temperature T_0, arising from the hopping or tunneling nature of the single quasiparticle c-axis propagation. Hence, the less-singular contribution *of the opposite sign* to the conductivity arising from the fluctuation decrease of the single quasiparticle density of states (DOS) at the Fermi energy ϵ_F becomes dominant above T_0. The competition between these contributions gives rise to a peak, or maximum in $\rho_c(T)$ just above T_c. The Maki–Thompson (MT) diagrams can be relevant, but such contributions were omitted in that treatment [8], as the assumption of strong pairbreaking was made.

In this work, we have studied the leading contributions to $\rho_c(T, H)$ arising from superconducting fluctuations of the order parameter in the presence of a perpendicular magnetic field $\mathbf{H}||\hat{c}$. Our results, which are valid for arbitrary impurity scattering, indicate that the peak in $\rho(T, H)$ increases with increasing H, and shifts to lower T values. They are then compared with recent experiments on YBCO and BSCCO. The details of this calculation will be presented elsewhere [9]. We use units in which $\hbar = k_B = c = 1$.

We assume free intralayer quasiparticle motion with effective mass m, Fermi velocity v_F, non-magnetic impurity scattering lifetime τ, effective pair-breaking lifetime τ_ϕ, interlayer hopping strength J, and c-axis repeat distance s. For this model [10], the normal single-spin quasiparticle density of states $N(0) = m/(2\pi s)$. For simplicity, we assume $J\tau << 1$. There are several functions of $\Lambda \equiv 1/(4\pi\tau T)$ which enter into the theory. The first of these is $\eta = \frac{1}{2}v_F^2\tau^2 f(\Lambda)$, where

$$f(\Lambda) = \psi(1/2) - \psi(1/2 + \Lambda) + \Lambda\psi'(1/2),$$

which is the positive constant which enters into the current expression in the phenomenological time-dependent Ginzburg–Landau (GL) theory in two dimensions, where $\psi'(x)$ is the derivative of the digamma function and v_F is the Fermi velocity for intralayer propagation. For $\tau T << 1$, η reduces to $\pi v_F^2\tau/(16T)$, and for $\tau T >> 1$, η approaches $7\zeta(3)v_F^2/(32\pi^2T^2)$, where $\zeta(x)$ is the Riemann zeta function.

There are also two functions $\kappa(\Lambda) = g(\Lambda)/[\pi^2 f(\Lambda)]$ and $\tilde{\kappa}(\Lambda) = h(\Lambda)/[\pi^2 f(\Lambda)]$, where

$$g(\Lambda) = \psi'(1/2 + \Lambda) - 2\Lambda\psi''(1/2),$$

$$h(\Lambda) = \psi'(1/2 + \Lambda) - \psi'(1/2) - \Lambda\psi''(1/2).$$

The parameter κ depends strongly upon τT, approaching the constant 0.691 for $\tau T << 1$, and behaving as $9.384(\tau T)^2$ for $\tau T >> 1$, whereas $\tilde{\kappa}$ is nearly constant, varying between 0.3455 for $\tau T << 1$ and 0.5865 for $\tau T >> 1$.

Second, we define $r = 4\eta J^2/v_F^2$, where $r(T_c) = 4\xi_\perp^2(0)/s^2$ is the usual anisotropy parameter [10] characterizing the dimensional crossover from the '2D' to the '3D' regimes in the thermodynamic fluctuation behavior at T_0 given by $\xi_\perp(T_0) = s/2$, and $\xi_\perp(0)$ is the GL coherence length in the c-axis direction at $T = 0$. The overall effect of pairbreaking is incorporated in the parameter $\gamma = 2\eta/[v_F^2\tau\tau_\phi]$. The magnetic induction B enters through the parameter $\beta = 4\eta eB$, where $\beta(T_c) = 4B\pi\xi_{||}^2(0)/\Phi_0$, $\Phi_0/[2\pi\xi_{||}^2(0)]$ is $H_{c2,\perp}(0)$ extrapolated from its slope near to T_c, Φ_0 is the flux quantum, and $\xi_{||}(0)$ is the GL coherence length parallel to the layers at $T = 0$. Near to $T_c(B)$, we set $B = H$.

The main (singular) temperature dependence of the various terms is incorporated in

$$\epsilon_B = \epsilon + \psi(1/2 + \beta/\pi^2) - \psi(1/2) \approx \epsilon + \beta/2,$$

or $\epsilon_B \approx [T - T_c(B)]/T_{c0}$, where $\epsilon = \ln(T/T_{c0}) \approx [T - T_{c0}]/T_{c0} << 1$.

We have evaluated the main (AL, DOS, and MT) contributions to the fluctuation conductivity in the presence of a magnetic field $\mathbf{H}||\hat{c}$, neglecting non-local magnetic field effects. The MT contribution contains a regular part, independent of τ_ϕ, and an anomalous part, which depends strongly upon τ_ϕ. We find for $J\tau \ll 1$,

$$\sigma_{zz}^{AL} = \frac{e^2 s r^2 \beta}{128\eta}\sum_{n=0}^{\infty}\frac{1}{[(\epsilon_B + \beta n)(\epsilon_B + \beta n + r)]^{3/2}}, \quad (1)$$

$$\sigma_{zz}^{DOS} = -\frac{e^2 sr\kappa\beta}{16\eta}\sum_{n=0}^{1/\beta}\frac{1}{[(\epsilon_B+\beta n)(\epsilon_B+\beta n+r)]^{1/2}}, \tag{2}$$

$$\sigma_{zz}^{MT(\mathrm{reg})} = -\frac{e^2 s\tilde{\kappa}\beta}{8\eta}\sum_{n=0}^{\infty}\left(\frac{\epsilon_B+\beta n+r/2}{[(\epsilon_B+\beta n)(\epsilon_B+\beta n+r)]^{1/2}}-1\right), \tag{3}$$

and

$$\sigma_{zz}^{MT(\mathrm{an})} = \frac{e^2 s\beta}{16\eta(\epsilon-\gamma)}\sum_{n=0}^{\infty}\left(\frac{\gamma_B+\beta n+r/2}{[(\gamma_B+\beta n)(\gamma_B+\beta n+r)]^{1/2}}-\frac{\epsilon_B+\beta n+r/2}{[(\epsilon_B+\beta n)(\epsilon_B+\beta n+r)]^{1/2}}\right), \tag{4}$$

where $\gamma_B \equiv \gamma+\beta/2$. We note that (1) was obtained previously [11].

For weak fields, we may use the Euler–Maclaurin approximation formula to expand (1)-(4) in powers of β. The fluctuation conductivity can then be shown to have a minimum (i. e., the resistivity has a maximum) at the temperature T_m, which is usually in the 2D regime. Letting $\epsilon_m = \ln[T_m/T_{c0}]$, in the 2D regime we have

$$\epsilon_m/r \approx \frac{1}{(8r\kappa)^{1/2}}\left(1-\frac{5\beta^2\kappa}{3r}\right)-\frac{\tilde{\kappa}}{8\kappa}+\frac{1}{16\gamma\kappa}, \tag{5}$$

which is satisfied for $r\kappa << 1$. The corrections due to the MT terms are usually small, and the magnetic field reduces T_m by an amount proportional to B^2.

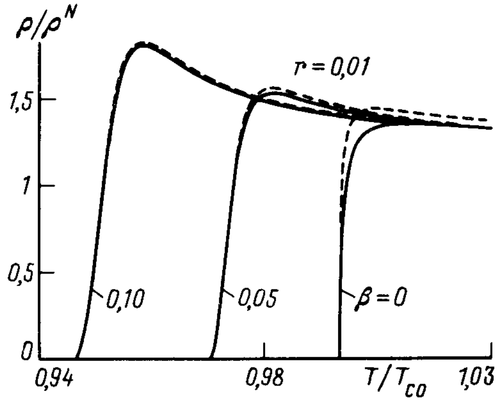
For strong fields, $\epsilon_B << \beta << 1$, the $n=0$ term in each sum in (1)–(4) dominates the behavior of each contribution. In this limit, we usually have $\gamma_B >> \max(\epsilon_B, r)$, since the combined zero-field and magnetic pairbreaking in $\gamma_B = \gamma+\beta/2$ is generally sufficiently strong that $r << \gamma_B$. As for the weak field case, the resistive maximum occurs at T_m obtained by setting $\epsilon_{Bm} = \epsilon_B(T_m)$. Again, T_m is usually in the 2D regime and given by

$$\epsilon_{Bm}/r \approx (3/8r\kappa)^{1/2} - \tilde{\kappa}/(2\kappa) + 1/(4\gamma_B\kappa). \tag{6}$$

Note that (6) is very similar to (5), the leading term differing by $\sqrt{3}$, and the zero-field correction terms by a factor of 4. However, since both ϵ_{Bm} and γ_B depend *linearly* upon β, the magnetic field decreases T_m by an amount linear in B.

The normal state c-axis conductivity $\sigma_{zz}^N = N(0)J^2\tau e^2 s^2/2$ for this model. The total conductivity $\sigma_{zz} = 1/\rho_{zz}$ is obtained by adding the normal state and fluctuation conductivities. In writing the c-axis resistivity ρ_{zz}, we therefore normalize our results to $\rho_{zz}^N = 1/\sigma_{zz}^N$. This introduces the Fermi energy $E_F = mv_F^2/2$. For the Boltzmann equation (and for our diagrammatic scattering procedure) to be valid, we must have $E_F\tau >> 1$. Furthermore, we must choose E_F/T_{c0} sufficiently large that $|\sigma_{zz}^{fl}/\sigma_{zz}^N| << 1$. For relatively clean materials, $\kappa >> 1$, so this requires a rather large value of E_F/T_{c0}.

In Figure we have plotted $\rho_{zz}(T,B)/\rho_{zz}^N$ for $r(T_{c0}) = 0.01$, using the full expressions (1)-(4), with a Gaussian cutoff in (2) for smoothness. We have chosen $\tau T_{c0} = 1$, which is close to that expected for the high-T_c cuprates. We have shown the behavior for two values of the pairbreaking parameter $\tau_\phi T_{c0}$, corresponding to strong ($\tau_\phi T_{c0} = 1$, dashed curves) and moderate ($\tau_\phi T_{c0} = 10$, solid curves) values. We have chosen $E_F/T_{c0} = 300$, so that ρ_{zz}/ρ_{zz}^N is not too much larger than unity at $T/T_{c0} \geq 1.03$. Decreasing E_F/T_{c0} with the other parameters



Plots of $\rho_{zz}(T,B)/\rho_{zz}^N$ versus T/T_{c0}, for $r(T_{c0}) = 0.01$, $\tau T_{c0} = 1$, $\tau_\phi T_{c0} = 1$ (dashed),10 (solid), $E_F/T_{c0} = 300$, and $\beta(T_{c0}) = 0$, 0.05, 0.10

fixed increases ρ_{zz}/ρ_{zz}^N, and enhances the magnitude of the peak. In addition to zero-field curves, curves for $\beta(T_{c0}) = 0.05$ and 0.10 are shown. As seen from these figures, decreasing $r(T_{c0})$ with the other parameters fixed changes the behavior dramatically. For $r(T_{c0}) = 0.1$, which is roughly appropriate for YBCO, there is essentially no peak for this range of parameters, consistent with experiments on that compound [2-4].

For BSCCO, we expect $r(T_{c0}) \approx 0.01$, as in Figure. As seen in this figure, there is a peak at all field values, which increases in magnitude with increasing field, as observed in the experiments [1,5]. Furthermore, curves with $\tau_\phi T_{c0} = 10$ appear more similar to the experiments [5] than do the $\tau_\phi T_{c0} = 1$ curves, since the experimental curves with increasing field strength lie on top of those for smaller field strengths, at least for temperatures above the maxima. The overall magnitude of the effect in Figure is less than observed in BSCCO, but it could be increased by decreasing the E_F/T_{c0} value. We expect that the renomalization of $T_c(B)$ arising from critical fluctuations will bring our results in more quantitative agreement with experiment. Hence, our results suggest that in BSCCO, the pairbreaking lifetime $\tau_\phi T_{c0} \approx 10$, which is *greater* than previously estimated for YBCO [12,13].

We remark that several of the Tl-based cuprates, such as $Tl_2Ba_2CaCu_2O_{8+\delta}$, are much more anisotropic as determined from torque measurements [14] than is BSCCO, resulting in $r(T_{c0})$ values in the range $0.001 - 0.0001$. Similar huge anisotropies were observed [14] in an organic layered superconductor, but those systems are probably in the dirty limit. In such highly anistropic materials, our theory predicts a sharp peak in the c-axis resistivity, with a magnetic field dependence that can be even more dramatic than that observed [5] in BSCCO. In addition, the peaks predicted are rather insensitive to the pairbreaking rate, except for very low field values. That is, the magnetic field produces sufficient pairbreaking to allow for a sharp peak.

In conclusion, it appears that superconducting fluctuations can account for the c-axis resitivity behavior observed in the high-T_c cuprates. In addition, it appears that the pairbreaking rate in the cuprates may be much less than previously thought. Further study is required in order to bring the theory into quantitative agreement with experiment.

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