Magnetic field enhancement of the c-axis resistivity peak near T_c in layered superconductors

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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, which lead to a peak in the overall c-axis resistivity $\rho_c(T)$ above T_c. In a magnetic field $\mathbf{H} \| \hat{\mathbf{c}}$, 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 many groups. (Refs. 2–4). Reproducible results on untwinned samples of $YBa_2Cu_3O_{7-\delta}$ (YBCO) (Refs. 2–4) and on $Bi_2Sr_2CaCu_2O_{8+\delta}$ (BSCCO) (Refs. 1 and 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 $\mathbf{H} \| \hat{\mathbf{c}}$ 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 underlying 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, and that it arises 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, which arises 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 a maximum in $\rho_c(T)$ just above T_c. The Maki–Thompson (MT) diagrams might be relevant, but such contributions were omitted in that treatment,[8] since strong pairbreaking was assumed to occur.

In this letter we report the results of an experimental study of the leading contributions to $\rho_c(T,H)$ which arise from superconducting fluctuations of the order parameter in the presence of a perpendicular magnetic field $\mathbf{H}\|\ \hat{\mathbf{c}}$. 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. These results 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 an effective mass m, Fermi velocity v_F, nonmagnetic 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 is $N(0)=m/(2\pi s)$. For simplicity, we assume $J\tau\ll 1$. There are several functions of $\Lambda\equiv 1/(4\pi\tau T)$ which are found in the theory. The first of these functions is $\eta=\frac{1}{2}v_F^2\tau^2 f(\Lambda)$, where

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

This is the positive constant which is contained in 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 of the intralayer propagation. For $\tau T\ll 1$, η reduces to $\pi v_F^2\tau/(16T)$, and for $\tau T\gg 1$, 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

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

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

The parameter κ depends strongly on τT. It approaches the constant 0.691 for $\tau T\ll 1$ and behaves as $9.384(\tau T)^2$ for $\tau T\gg 1$, whereas $\tilde{\kappa}$ is nearly constant; it varies between 0.3455 for $\tau T\ll 1$ and 0.5865 for $\tau T\gg 1$.

Second, we define $\tau=4\eta J^2/v_F^2$, where $\tau(T_c)=4\xi_\perp^2(0)/s^2$ is the usual anisotropy parameter[10] which characterizes 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 pair breaking 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 T_c, Φ_0 is the flux quantum, and $\xi_\perp(0)$ is the GL coherence length which is parallel to the layers at $T=0$. Near $T_c(B)$, 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 in $\epsilon_\beta \approx [T - T_c(B)]/T_{c0}$, where $\epsilon = \ln(T/T_{c0}) \approx [T - T_{c0}]/T_{c0} \ll 1$.

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

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

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

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

and

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

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. Setting $\epsilon_m = \ln[T_m/T_{c0}]$, in the 2D regime we have

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

which is satisfied for $\tau\kappa \ll 1$. 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, $\epsilon_B \ll \beta \ll 1$, the $n=0$ term in each sum in (1)–(4) dominates the behavior of each contribution. In this limit, we usually have $\gamma_B \gg \max(\epsilon_B, \tau)$, since the combined zero-field and magnetic pair breaking in $\gamma_B = \gamma + \beta/2$ is generally sufficiently strong that $\tau \ll \gamma_B$. As for the case of the weak field, 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 is given by

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

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

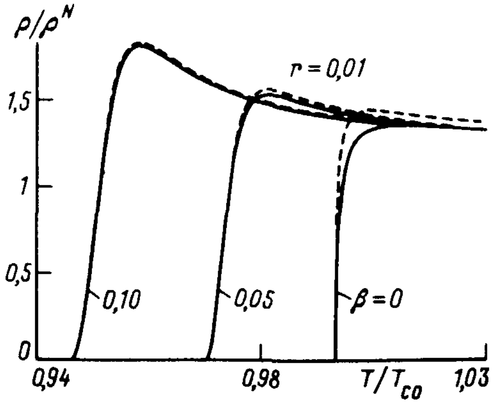

FIG. 1. Plots of $\rho_{zz}(T,B)/\rho_{zz}^N$ versus T/T_{c0}, for $\tau(T_{c0})=0.01$, $\tau T_{c0}=1$, $\tau_\phi T_{c0}=1$ (dashed), 10 (solid), $E_F/T_{c0}=300$, and $\beta(T_{c0})=0$, 0.05, 0.10.

The normal state c-axis conductivity is $\sigma_{zz}^N = N(0)J^2\tau e^2 s^2/2$ for this model. The total conductivity $\sigma_{zz}=1/\rho_{zz}$ is obtained by adding the normal state and the 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 $E_F\tau \gg 1$. Furthermore, we must choose E_F/T_{c0} sufficiently large in order that $|\sigma_{zz}^{fl}/\sigma_{zz}^N| \ll 1$. Since for relatively clean materials we have $\kappa \gg 1$, this requires a rather large value of E_F/T_{c0}.

In Fig. 1 we have plotted $\rho_{zz}(T,B)/\rho_{zz}^N$ for $\tau(T_{c0})=0.01$, using 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 pair-breaking parameter $\tau_\phi T_{c0}$, which corresponds 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 $T/T_{c0}\geqslant 1.03$. Decreasing E_F/T_{c0}, while holding the other parameters constant, 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 in Fig. 1. As can be seen from these curves, decreasing $\tau(T_{c0})$, while holding the other parameters fixed, changes the behavior dramatically. For $\tau(T_{c0})=0.1$, 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 $\tau(T_{c0})\approx 0.01$, as in Fig. 1. As can be 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 $\tau_\phi T_{c0}=10$ are more similar to the experiments[5] than are 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 total strength of the effect in Fig. 1 is less than that observed in BSCCO, but it could be increased by decreasing the E_F/T_{c0} value. We expect that the renormalization of $T_c(B)$, which arises from critical fluctuations, will bring our results in better quantitative agreement with experiment. Hence, our results suggest that in BSCCO, the pair-breaking lifetime is $\tau_\phi T_{c0} \approx 10$, which is *greater* than previously estimated for YBCO.[12,13]

We note 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, which accounts for the $\tau(T_{c0})$ 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 anisotropic 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, which were predicted are rather insensitive to the pair-breaking rate, except for very low field values. In other words, the magnetic field produces sufficient pair breaking to allow for a sharp peak.

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

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