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Experimental evidence in favor of the fluctuation origin of the transverse-resistance increase near the edge of the superconducting transition in $Bi_2Sr_2CaCu_2O_{8+\delta}$

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The experimentally observed increase of the transversal resistance near the edge of the transition in $Bi_2Sr_2CaCu_2O_{8+\delta}$ epitaxial film is analyzed on the basis of a recent theoretical proposal on its fluctuation origin [L. Ioffe, A. I. Larkin, A. A. Varlamov, and L. Yu, Phys. Rev. B (to be published). It is shown that the excess resistance in a wide range of temperatures increases logarithmically as the superconducting transition temperature is approached from above, in agreement with the theoretical consideration. These results are also compared with the available data on $YBa_2Cu_3O_7$ compounds.

I. INTRODUCTION

The maximum in the temperature dependence of the resistivity of the high-T_c superconductors near the edge of a transition has a long history. This type of semiconductor behavior was often observed in ceramic samples. As the quality of samples has been improving and single crystals are becoming available, the shape of the transition for in-layer measurements has appeared to be more and more traditional,[1] i.e., a linear temperature dependence in the normal state and a sharp drop at T_c. At the present time, the quality of samples permits one to measure reliably the conductivity not only in the *ab* plane, but in the *c* direction as well. As a result, the increase of the transversal resistance was observed unambiguously for $Bi_2Sr_2CaCu_2O_8$ compounds.[2,3] The situation is less clear for $YBa_2Cu_3O_7$ compounds where the temperature variation of the transverse resistance is strongly sample dependent.

As remarked by Anderson,[4] the difference in the temperature dependence of the longitudinal and transverse resistivity is the most difficult point to explain in the framework of conventional Fermi-liquid theory. Recently, a possible interpretation of this difference was identified with a striking dissimilarity in the role of superconducting fluctuations in the *ab* plane and *c*-axis direction.[5] It turns out that the usual paraconductivity in the case of transversal current shows up only in the second order of the interlayer transparency (fourth power of the transverse hopping integral). On the other hand, the less singular, as a function of the reduced temperature, negative contribution, originating from the decrease of the single-electron density of states in each layer due to the presence of fluctuation pairing, is proportional to the first order of transparency and becomes more important. It gives rise to a logarithmic increase of the transverse resistivity, which can be observed in a sufficiently wide range of temperatures. As a consequence, these two competing contributions of opposite sign with very different temperature dependence lead to the appearance of a maximum in $\rho(T)$ in the immediate vicinity of the critical temperature.

In this paper we first report our experimental results on the transverse resistance of $Bi_2Sr_2CaCu_2O_{8+\delta}$ epitaxial films and we show that they cannot be fitted by a simple power-law temperature dependence in a sufficiently wide range or by a more sophisticated empirical formula[6,7] below 150 K (Sec. II). Then we briefly outline the theoretical results of Ref. 5 in a form more convenient for comparison with experiments and more likely to make this paper self-contained (Sec. III). Furthermore, we fit the data with the theoretical formulas of Ref. 5 and the agreement turns out to be very good in a wide range of temperatures above T_c (Sec. IV). We consider this as strong evidence in favor of the fluctuation origin of the resistance increase near T_c. Finally, we make some concluding remarks and comments on the difference between $YBa_2Cu_3O_7$ and $Bi_2Sr_2CaCu_2O_8$ (Sec. V).

II. EXPERIMENT

The investigation of the fluctuation conductivity in high-T_c superconductors requires careful measurements of the resistivity behavior above the transition temperature. In the case of in-plane fluctuation conductivity in $Bi_2Sr_2CaCu_2O_8$ systems, the experiments indicate a two-dimensional behavior of fluctuations in a wide range of temperatures $-4<\ln[(T-T_c)/T_c]<-2$ (see Ref. 8).

On the other hand, the experimental investigation of the fluctuation conductivity along the *c* axis has to be carried out either on single crystals or monocrystalline films. Unfortunately, to our knowledge, no one has succeeded so far in growing truly monocrystalline films of $Bi_2Sr_2CaCu_2O_8$ with the c axis in plane. Moreover, single crystals of $Bi_2Sr_2CaCu_2O_8$, grown by the usual flux method, contain a large number of extended defects perpendicular to the c direction. To overcome, as much as possible, this difficulty, single crystals are usually cleaved into extremely thin ab-plane sheets. Therefore measurements of $\rho_c(T)$ have to be carried out along the c crystallographic direction on ab-oriented single crystal sheets of a few μm thick only. Such a setup gives rise to an unavoidable degradation of the ρ_c measurements which can hinder the small fluctuation effects on the resistivity itself.

To avoid this problem, we have measured the resistivity of monocrystalline films grown by liquid-phase epitaxy (LPE) on slightly ab-misoriented $NdGaO_3$ substrates. LPE has proved to be a suitable technique to grow monocrystalline $Bi_2Sr_2CaCu_2O_8$ films truly epitaxial relative to the substrate.[9–11] The nice feature of the LPE films is the very narrow mosaic spread (less than 0.1°). The structural and transport characterizations of the $Bi_2Sr_2CaCu_2O_8$ films have been discussed in Ref. 12. The crystal $NdGaO_3$, which has an orthorhombic structure with $a=5.42$ Å, $b=5.50$ Å, and $c=7.71$ Å, turned out to be a suitable substrate for the $Bi_2Sr_2CaCu_2O_8$ films (with lattice parameters $a=5.41$ Å, $b=5.43$ Å, and $c=30.81$ Å). The (001)-oriented $NdGaO_3$ slices provide a satisfactory in-plane match between the a and b lattice parameters of the film and substrate and, therefore, can be used to grow ab-oriented $Bi_2Sr_2CaCu_2O_8$ films.

We have grown $Bi_2Sr_2CaCu_2O_8$ films on $NdGaO_3$ substrates slightly misoriented relative to the (001) plane. The misorientation angle Φ (the angle between the normal to the physical surface of the film and the c-crystallographic axis) was between 1° and 3°, i.e., much larger than the mosaic spread of the epitaxial $Bi_2Sr_2CaCu_2O_8$ films. Nevertheless, Φ values of a few degrees will still allow the growth of good quality epitaxial films. In Fig. 1 the behavior of resistivity vs temperature is shown for a film grown on a $NdGaO_3$ substrate with $\Phi=2.5°$. For comparison, the behavior of resistivity vs temperature for a typical film grown on a well-oriented substrate (with Φ much smaller than the mosaic spread) is also reported in the same figure. The resistivity for the film grown on a misoriented substrate shows a nonmetallic behavior with a maximum close to the transition temperature. The film thickness is 1 μm.

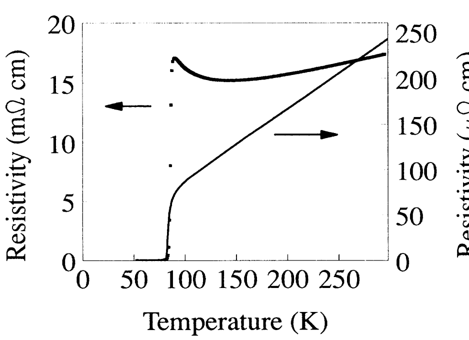

FIG. 1. Resistivity vs temperature curve for an epitaxial $Bi_2Sr_2CaCu_2O_8$ film grown on a 2.5°-misoriented $NdGaO_3$ substrate (squares). For comparison, the resistivity vs temperature curve of a $Bi_2Sr_2CaCu_2O_8$ film grown on a well-oriented $NdGaO_3$ substrate is also reported (solid line).

The resistivity measurements were carried out by the standard four-probe technique on a narrow strip 50 μm wide and 1 mm long. The strip was obtained by the standard photolithographic technique. The misorientation between the strip and ab plane was 2.5°. Since the ratio of $\rho_c/\rho_{ab}\approx 10^4$, the resistance of the strip, even for a misorientation as small as 2.5°, is still dominated by the contribution along the c axis. As a matter of fact, the resistivity shown in Fig. 1 just above the transition temperature is larger than that found for films grown on well-oriented substrates (with a misorientation smaller than the mosaic spread of the film) by a factor of about 300. A straightforward estimate of ρ_c from the data of Fig. 1 gives $\rho_c(90\ \mathrm{K})\sim 1\ \Omega\,\mathrm{cm}$, in good agreement with resistivity values reported from measurements on single crystals. All these considerations imply that the contribution from the in-plane conduction in Fig. 1 is negligible.

Various models have been proposed to explain the unusual nonmetallic behavior of the out-of-plane conduction of $Bi_2Sr_{3-x}Ca_xCu_2O_{8+\delta}$ compounds, including interplanar tunneling[13] and percolation conduction enhanced by defect scattering.[14] A simple empirical expression was first proposed to explain the behavior of the electrical conductivity along the c axis in the case of $YBa_2Cu_3O_{7-\delta}$ (Ref. 6) and then extended to $Bi_2Sr_{3-x}Ca_xCu_2O_{8+\delta}$ (Ref. 7):

$$\rho_c(T)=\rho_0 T^{\alpha}\exp\left[\frac{\Delta}{k_B T}\right], \tag{1}$$

where α is about 0.7 and Δ represents a small activation energy of the order of 10 meV. However, attempts to fit $\rho_c(T)$ with the above expression in the whole range from room temperature to the transition temperature were not fully satisfactory. Below about 150 K the experimental curve of $\rho_c(T)$ deviated from the empirical expression (1). The deviation was particularly evident between 100 K and the transition temperature where the experimental $\rho_c(T)$ greatly overshoots the empirical formula. It was also noted in Ref. 7 that "precisely in this temperature range the ab-plane electrical resistivity deviates from a purely linear temperature dependence." An attempt to fit our experimental data using the empirical formula (1) gave quite poor results similar to those described in Ref. 7.

III. THEORY

As mentioned in the Introduction, there has been a recent proposal[5] to interpret the increase of the transverse resistance near the transition temperature in terms of superconducting fluctuations. Here we summarize the main results of this theoretical consideration.

In order to calculate the transverse component of the fluctuation conductivity in a layered superconductor, we adopt the electron spectrum in the form of a modulated cylinder:

$$\xi(\mathbf{p})=\varepsilon(\mathbf{p})-\varepsilon_F=v_F(|p_{\parallel}|-p_F)+w\cos(p_{\perp}a)\ , \qquad (2)$$

where $\mathbf{p}=(\mathbf{p}_{\parallel},p_{\perp})$ is the electron momentum, w the electron hopping integral between layers, and a the interlayer distance.

The transverse fluctuation conductivity may be easily expressed in terms of the analytically continued operator of the electromagnetic response $Q^{\mathrm{fl}}_{\alpha\beta}(\omega)$:[15]

$$\sigma_{\perp}^{\mathrm{fl}}=\lim_{\omega\to 0}\frac{[Q_{\perp}^{\mathrm{fl}}(\omega)]^R}{-i\omega}\ . \qquad (3)$$

The diagrams for $Q^{\mathrm{fl}}_{\alpha\beta}$ are presented in Fig. 2. We consider the case of a clean metal ($T\tau\gg 1$), and in the calculation of paraconductivity, we neglect the electron scattering described by the relaxation time τ. We focus on the most interesting case of low transparency $w\ll T_c$, which is materialized for $Bi_2Sr_2CaCu_2O_8$ compounds and gives rise to an unusual competition between different contributions and a crossover at temperatures $T-T_c\ll T_c$ for the ab in-plane fluctuation conductivity.[8] However, the hierarchy of the fluctuation contributions in the c direction differs dramatically from the in-layer situation and leads to a nontrivial increase of resistivity at the edge of transition in the c direction, as we will see below.

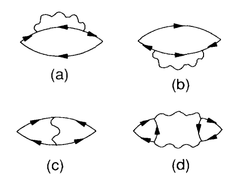

FIG. 2. Diagrams for the different fluctuation contributions in the operator of the electromagnetic response.

A. Transversal paraconductivity

The calculation of the Aslamazov-Larkin (AL) contribution [Fig. 2(d)] is performed in the usual way[15,8] with velocities

$$v_{\perp}(p)=\frac{\partial\varepsilon}{\partial p_{\perp}}=-wa\sin(p_{\perp}a)\ , \qquad (4)$$

in each external vertex of the diagram. The fluctuation propagator (the wavy line in Fig. 2) for layered superconductors has been discussed repeatedly (see, for example, Refs. 8 and 16):

$$L^R(\mathbf{q},\omega)=-\frac{1}{\rho}\frac{1}{[(T-T_c)/T_c]-i\pi\omega/8T_c+\frac{3}{2}\eta q_{\parallel}^2+6\eta(w^2/v_F^2)\sin^2(q_{\perp}a/2)}\ , \qquad (5)$$

where ρ is the density of states and $\eta=7\zeta(3)v_F^2/48\pi^2T_c^2$ is the well-known parameter of the Ginzburg-Landau theory in the clean case.

After a straightforward calculation of the diagram Fig. 2(d), one finds the general expression for the AL contribution to the transverse conductivity when $T-T_c\ll T_c$:

$$\sigma_{\mathrm{AL}}^{\perp}=\frac{e^2}{16a}\left[\frac{wa}{v_F}\right]^2\left\{\left[\left[\delta_0^2+\frac{T-T_c}{T_c}\right]^{1/2}-\left[\frac{T-T_c}{T_c}\right]^{1/2}\right]^2\Big/\delta_0^2\left[\frac{T-T_c}{T_c}\right]^{1/2}\left[\delta_0^2+\frac{T-T_c}{T_c}\right]^{1/2}\right\}\ , \qquad (6)$$

where $\delta_0^2=7\zeta(3)w^2/8\pi^2T_c^2$ is the parameter characterizing the effective transverse size of the Cooper pair in the c direction,

$$\xi_{\perp}(T)=\frac{\delta_0 a}{[(T-T_c)/T_c]^{1/2}}\ , \qquad (7)$$

and its ratio to the lattice constant a determines the dimensionality of the fluctuating Cooper pairs [at the crossover temperature $T_{\mathrm{cr}}-T_c\sim w^2/T_c$, the condition $\xi_{\perp}(T_{\mathrm{cr}})\sim a$ is fulfilled]. One can see from (6) that in the immediate vicinity of the transition, when $\xi_{\perp}(T)\gg a$ ($T-T_c\ll w^2/T_c$),

$$\sigma_{\mathrm{AL}}^{\perp(3\mathrm{D})}=\frac{e^2}{8a}\left[\frac{2\pi^2}{7\zeta(3)}\right]^{1/2}\left[\frac{wa}{v_F}\right]\left[\frac{T_ca}{v_F}\right]\left[\frac{T_c}{T-T_c}\right]^{1/2}, \qquad (8)$$

and the behavior of the paraconductivity is evidently three dimensional, because the size of the Cooper pair is much larger than the interlayer distance.

At temperatures higher than the crossover point ($T-T_c\geq w^2/T_c$) $\xi_{\perp}(T)\leq a$, one finds from (6) a quite surprising result:

$$\sigma_{\mathrm{AL}}^{\perp(0\mathrm{D})}=\frac{7\zeta(3)e^2}{2^9\pi^2a}\left[\frac{wa}{v_F}\right]^2\left[\frac{w}{T_c}\right]^2\left[\frac{T_c}{T-T_c}\right]^2\ . \qquad (9)$$

First of all, the critical exponent 2 in this range of temperature follows zero- (not two-) dimensional behavior of fluctuation conductivity (as, for example, in granular superconductors). Of course, the fluctuating Cooper pairs are two dimensional (2D) and contribute as 2D objects to the fluctuation conductivity $\sigma_{\mathrm{AL}}^{\parallel}$ in the ab planes.[8,16] However, their motion in the c direction has a hopping character, so due to the necessity of replacement in the next layer of both electrons during the Cooper pair lifetime $\Delta t\sim 1/(T-T_c)$ the appropriate charge transfer becomes possible in the second order of the barrier transparency only and the critical exponent changes from 1 to 2. This is the fact that reflects the effective zero-dimensionality of the Cooper-pair motion in the c direction.

Another important feature of (9) is that $\sigma_{AL}^{\perp(0D)}$ turns out to be proportional to w^4, i.e., the second order of the interlayer transparency in the language of the Lawrence-Doniach model.[17] It means that besides the natural anisotropy of the obtained result, an additional small parameter δ_0^2 is contained in it. As we will see below, this feature changes drastically the usual hierarchy of the fluctuation contributions. Contrary to the in-plane case, the contributions from the first three diagrams in Fig. 2 have to be taken into account for considering transverse conductivity.

B. Particular role of density-of-states fluctuations

In the case of $\sigma_{fl}^{\parallel}$, the contributions from the first three diagrams cancel each other *exactly* for dc fluctuation conductivity.[8] In contrast to the in-plane situation, in the case under consideration such a cancellation is no longer valid. In fact, the AL contribution is proportional to w^4. One can easily see that the contribution of the Maki-Thompson- (MT)-type diagram [Fig. 2(c)] gives nonvanishing results only in the fourth order of w, too (but it is less singular as a function of the reduced temperature in comparison with $\sigma_{AL}^{\perp}$). On the contrary, the contribution of the first two diagrams turns out to be proportional to w^2 and hence in the problem under discussion it will be of primary importance.

As discussed earlier,[18,19] these two diagrams describe the effect due to the change of the single-electron density of states. In fact, some electrons in each layer are involved in fluctuation Cooper pairing in every moment, and hence the number of electrons participating in charge transport (tunneling current) in the c direction effectively decreases.

In the above discussion of paraconductivity, there was no need to discuss electron scattering. Now, considering single-electron conductivity, we have to take into account some scattering mechanisms because they set the limit of the conductivity itself. The origin of electron scattering was discussed in detail in Ref. 5, and now we adopt the simplest case $T\tau \geq 1$ (but not very large) and evaluate the expression for single-electron conductivity,

$$\sigma_{\perp}(\omega) = -2e^2 \int \frac{d^3p}{(2\pi)^3} \int_{-\infty}^{\infty} d\varepsilon\, v_{\perp}^2 \frac{\partial n(\varepsilon)}{\partial \varepsilon} \,\mathrm{Im} G^A(\mathbf{p},\varepsilon+\omega) \times \mathrm{Im} G^A(\mathbf{p},\varepsilon) , \quad (10)$$

where $n(\varepsilon)$ is the Fermi distribution.

The advanced Green's function $G^A(\mathbf{p},\varepsilon)$ has to take into account both electron scattering (to all orders of the perturbation theory) and the possibility of fluctuation pairing of electrons (to first order only). It is convenient to express it in terms of the self-energy operator

$$G^A(\mathbf{p},\varepsilon) = [G_0^A(\mathbf{p},\varepsilon)]^2 \Sigma^A(\mathbf{p},\varepsilon) , \quad (11)$$

where

$$G_0^A(\mathbf{p},\varepsilon) = \left[\varepsilon - \xi(\mathbf{p}) - \frac{i}{2\tau}\right]^{-1} \quad (12)$$

and

$$\Sigma^A(\mathbf{p},\varepsilon_n) = T \sum_{\Omega_\kappa} \int \frac{d\mathbf{q}}{(2\pi)^3} L(\mathbf{q},\Omega_\kappa) G_0(\mathbf{q}-\mathbf{p},\Omega_\kappa-\varepsilon_n) . \quad (13)$$

Accomplishing the analytical continuation of (13) and carrying out the integration over frequencies, one can find[5] for the density of states (DS) contribution [for the two-dimensional region $w^2/T_c \leq T - T_c \ll T_c$ and to leading order of $T_c/(T-T_c)$ only]

$$\sigma_{DS}^{\perp(2D)} = -\frac{\pi^2 e^2}{7\zeta(3)a}(T_c\tau)^2 \left[\frac{wa}{v_F}\right]^2 \ln\left[\frac{T_c}{T-T_c}\right] . \quad (14)$$

The total fluctuation correction in transverse conductivity for temperatures $w^2/T_c \leq T - T_c \ll T_c$ is determined by the sum of contributions (6) and (14).

Let us now discuss the results obtained. First of all, the hopping character of the electron motion in the c direction changes drastically the usual hierarchy of the fluctuation corrections and, what is even more important, it changes the sign of the correction in a wide range of temperatures (if $w \ll T_c$). Indeed, in spite of the dramatic difference between the weak temperature-dependent $\sigma_{DS}^{\perp(2D)}$ $(\sim \ln[T_c/(T-T_c)])$ and the strongly temperature-dependent $\sigma_{AL}^{\perp(0D)}$ $(\sim [T_c/(T-T_c)]^2)$, the presence of an additional small transparency coefficient $(\sim w^2)$ in the latter gives rise to the domination of the 2D contribution in a sufficiently wide range of temperatures above T_c. The competition of these two contributions leads to the appearance of a maximum in the temperature dependence of the resistance at

$$T_m^{(2D)} = T_c \left[1 + \frac{7\zeta(3)}{32\pi^2} \frac{w}{T_c} \frac{\sqrt{2}}{T_c\tau}\right] \quad (15)$$

(of course, if this point belongs to the range of temperatures where both formulas are valid) and the resistivity decreases as the temperature is lowered further. At $T_{cr} - T_c \sim w^2/T_c$, the crossover from zero-dimensional behavior [Eq. (9)] to three-dimensional behavior [Eq. (8)] occurs. At this temperature $\xi_{\perp}(T_{cr}) \sim a$ and in some sense the barriers between the metallic layers disappear. Here $\sigma_{AL}^{\perp(3D)}$ does not contain w^4 (but w only). The effective size of Cooper pairs exceeds the interlayer distance, and at these temperatures the only scale left in the system is $\xi_{\perp}$.

One can easily check that, because of the presence of a small numerical coefficient in (15), in the case of not very small w and even not very large $T\tau$ the position of maximum can occur at a temperature below the crossover point $T_{cr} \sim w^2/T_c$. In this case expression (14) is already not valid, and in order to find the position of maximum one has to use the asymptotical expression (8) and evaluate expression (13) for the case of three-dimensional fluctuations. Simple analysis of formula (13) shows that at temperatures below the crossover point $\sigma_{DS}^{\perp(3D)}$ only weakly depends on the reduced temperature and is mainly determined by its value at the crossover point

$$\sigma_{DS}^{\perp(3D)} = -\frac{2\pi^2 e^2}{7\zeta(3)a}(T_c\tau)^2 \left[\frac{wa}{v_F}\right]^2 \ln\left[\frac{T_c}{w}\right] . \quad (16)$$

Then, for the position of the maximum in this case, one obtains the estimate

$$T_m^{(3D)} = T_c \left[1 + \frac{7\zeta(3)}{2^6\pi^2} \frac{1}{(T_c\tau)^2} \right] . \tag{17}$$

IV. COMPARISON OF THEORY WITH EXPERIMENTS

Now we try to interpret the experimental data of $\rho_c(T)$ reported in Sec. II in terms of the superconducting fluctuation model proposed by Ioffe *et al.*[5] The agreement between the experimental data for $\rho_c(T)$ in the temperature range from 120 K until the nearest vicinity of the critical temperature with the behavior predicted by the fluctuation model turns out to be very good.

In order to fit the experimental data with the results of fluctuation theory, we have to subtract the background of the normal-state resistance in the c direction, $\rho_l(T)$. As far as we know, the precise temperature dependence of the c-axis resistivity is still an open question. Some data show a more or less linear temperature dependence over certain range (see, e.g., Ref. 7), whereas some other data show a different dependence above the transition temperature (see, e.g., Refs. 2 and 3). Our data show a rather good linear temperature dependence over a wide range starting from room temperature. Although we do not yet fully understand the origin of such a dependence of resistance in the c direction, we would assume it as the background. For this purpose $\rho_l(T)$ has been estimated by extrapolating the experimental data from room temperature. The excess resistivity $\rho_{fl}(T) = \rho(T) - \rho_l(T)$ has then been fitted in the suitable temperature range $93 \leq T \leq 110$ K using expression (14). In accordance with the above theoretical discussion, the paraconducting contribution is expected to be negligible because of the small value of the hopping integral w. In Fig. 3 the resistivity behavior between 160 and 80 K is shown on an expanded scale for the same sample of Fig. 1. The theoretical expression (14) used to fit the experimental data contains only one free parameter, namely, the prefactor of $\ln[T_c/(T-T_c)]$. The critical temperature T_c used by us in (14) has been determined independently according to the usual criterium $\rho(T_c) = \rho_l(T_c)/2$. We would like to emphasize that our fitting was done under the assumption of a linear temperature dependence of the normal-state resistivity, but a more complicated expression, as given in (1), would obviously improve the quality of the fitting. Nevertheless, we prefer the simplest way to test the experimental data with theoretical predictions.

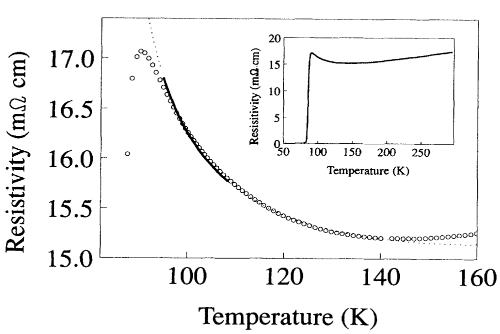

FIG. 3. Enlarged view of the resistivity vs temperature curve of the sample shown in Fig. 1. The solid line is the fit to the theoretical expression (14), while the dashed line shows the behavior of the fitting function outside the fitting region. In the inset the full resistivity vs temperature curve is plotted.

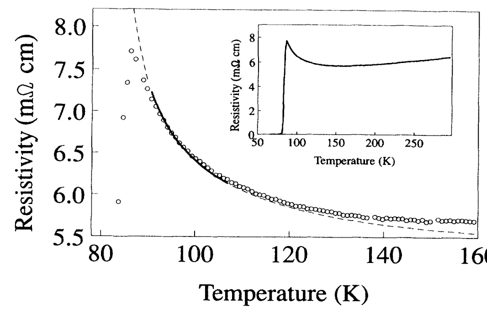

FIG. 4. Enlarged view of the resistivity vs temperature curve of an epitaxial $Bi_2Sr_2CaCu_2O_8$ film grown on a 1.3°-misoriented $NdGaO_3$ substrate. The solid line is the fit to the theoretical expression (14), while the dashed line shows the behavior of the fitting function outside the fitting region. In the inset the full resistivity vs temperature curve is plotted.

In Fig. 3 the solid line through the experimental points represents the fit with (14). For temperatures lower than 93 K, expression (14) overshoots the experimental value of $\rho(T)$. This effect can be attributed to the increase of the positive contribution from the paraconductivity [Eq. (6)]. At about 86 K, the competition of these two contributions determines the appearance of a maximum in $\rho(T)$. Then, as the temperature goes further down to T_c, the resistivity decreases, in a qualitative agreement with the theoretical prediction [Eq. (6)]. Unfortunately, the small number of our experimental data points does not permit us to fit the experimental results with the theoretical prediction in this region of temperatures.

For temperatures higher than 105–110 K, the results of the theory summarized above are not expected to be valid because of the assumption $T - T_c \ll T_c$ adopted above. However, the fit remains satisfactory up to about 140 K. This fact can be explained by the extremely slow logarithmic temperature dependence of the fluctuation correction.

The same measurements were performed on a second film grown on a substrate with $\Phi = 1.3°$. The results obtained for this film are shown in Fig. 4. Again, the fitting is very good. The resistivity is lower than that of the first film, as expected from the lower value of Φ.

V. DISCUSSION

Let us begin our discussion with an attempt to reconcile the position of the maximum in the experimental curve with that of the theoretical prediction. This procedure depends strongly on an assumption about the values of w and $T\tau$. The value of w for $Bi_2Sr_2CaCu_2O_8$ was estimated before from the crossover temperature[8] and from the anisotropy of the normal-state resistance as $\omega \sim 15$–20 K. $T\tau \gtrsim 1$ was adopted above. Hence the crossover temperature $T_{cr} - T_c \sim w^2/T_c \sim 4$–$5$ K and the maximum in the temperature dependence of the transverse resistance occurs exactly in the same temperature range. The estimation of the $\Delta T_m = T_m - T_c$ from formula (15) with $w \sim 15$ K and $T\tau$ even of order of 1 gives $\Delta T_m^{(2D)} \sim 0.5$ K, which is almost one order of magnitude less than the experimental value ($\Delta T_m^{expt} = 3$ K). More important is that $T_m^{(2D)}$ turns out to be below the crossover temperature; hence, for the estimation of T_m, we have to use the three-dimensional result $T_m^{(3D)}$ [Eq. (17)], which gives much better estimation of $\Delta T_m \sim 2$ K (at the same assumption $T\tau \sim 1$).

Therefore we have to adopt that in our experiments on $Bi_2Sr_2CaCu_2O_8$ films the crossover between 0D and 3D behavior of paraconductivity and the maximum in the transverse resistance occur in the same range of temperatures (a few degrees above the critical temperature). The intermediate value of the parameter $T\tau \sim 1$ is more or less appropriate, as one would expect.

It is worthwhile to mention that in the recent experiments of Baraduc *et al.*[20] on $YBa_2Cu_3O_7$ single crystals the crossover from 0D to 3D behavior has been observed. The authors of Ref. 20, independently of Ref. 5, found theoretically this unusual crossover behavior using the Ginzburg-Landau approach and have observed it experimentally. However, they did not find the maximum in the temperature dependence of the transverse resistance for $YBa_2Cu_3O_7$ samples. A plausible explanation of this discrepancy may come from the fact that the value of w for $YBa_2Cu_3O_7$ is larger compared with $Bi_2Sr_2CaCu_2O_8$; so paraconductivity (in spite of its proportionality to w^4) still dominates the contribution originating from density-of-states fluctuations in the entire region $T - T_c \ll T_c$. One may note in this connection that some $YBa_2Cu_3O_7$ samples did show some not very pronounced increase of the transverse resistance[21] which may be associated with the dependence of the hopping integral on the oxygenation state of the sample. Hess *et al.*[3] have measured the c-direction resistance on $Bi_2Sr_2CaCu_2O_8$ single crystals in the presence and absence of a magnetic field in the c direction. The transverse-resistance increase near T_c is very similar to our data, but a detailed comparison has to be made.

Finally, let us summarize the main results of this paper. The experimentally observed increase of the transverse resistance for $Bi_2Sr_2CaCu_2O_8$ epitaxial films in a wide range of temperatures is found to be in good agreement with the recently proposed fluctuation theory of this phenomenon.[5] Analysis of the experimental curve allows us to conclude that the values of the hopping integral and the elastic electron-scattering time are in agreement with the corresponding values extracted from in-plane fluctuation conductivity measurements and the anisotropy of normal-state conductivity.

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