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Crossover between Aslamazov-Larkin and short-wavelength fluctuation regimes in high-temperature-superconductor conductivity experiments

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We present paraconductivity measurements in three different high-temperature superconductors: a melt-textured $YBa_2Cu_3O_7$ sample, a $Bi_2Sr_2CaCu_2O_8$ epitaxial thick film, and a highly textured $Bi_2Sr_2Ca_2Cu_3O_{10}$ tape. The crossovers between different temperature regimes in excess conductivity have been analyzed. The Lawrence-Doniach (LD) crossover, which separates the two-dimensional and three-dimensional regimes, shifts from lower to higher temperatures as the compound anisotropy decreases. Once the LD crossover is overcome, the fluctuation conductivity of the three compounds shows the same universal behavior: for $\epsilon=\ln T/T_c>0.23$ all the curves bend down according to the $1/\epsilon^3$ law. This asymptotic behavior was theoretically predicted previously for the high-temperature region where the short-wavelength fluctuations become important. [S0163-1829(97)52122-9]

It is well known that, owing to strong anisotropy, high critical temperature, and low charge carrier concentration, thermodynamic fluctuations play an important role in the explanation of the normal-state properties of high-temperature superconductors (HTS). Just after the realization of high-quality epitaxial single-crystal samples, the in-plane fluctuation conductivity was investigated in detail and the Lawrence-Doniach (LD) crossover[1] between three-dimensional and two-dimensional regimes (or at least a tendency to it) was observed in the vicinity of T_c in the majority of HTS compounds.[2] Analogous phenomena were found in magnetic susceptibility, thermoconductivity,[3] and other properties of HTS.

How the LD crossover appears can be shown explicitly considering the model of an open electron Fermi surface which, for instance, can be chosen in the form of a ''corrugated cylinder.''[4] In this case the energy spectrum has the form

$$\xi(\mathbf{p})=\epsilon_0(\mathbf{p})+J\cos(p_\perp s)-E_F, \qquad (1)$$

where $\epsilon_0(\mathbf{p})=\mathbf{p}^2/(2m)$, $p\equiv(\mathbf{p},p_\perp)$, $\mathbf{p}\equiv(p_x,p_y)$ is a two-dimensional, intralayer wave vector, and J is an effective hopping energy. The Fermi surface is defined by the condition $\xi(p_F)=0$ and E_F is the Fermi energy. This spectrum is the most appropriate for strongly anisotropic layered materials where $J/E_F\ll 1$.

In the framework of the Ginzburg-Landau (GL) theory for an isotropic spectrum, the fluctuation contribution to the free energy of a superconductor slightly above the critical temperature can be presented as the sum over long-wavelength fluctuations.[5] This result can be easily extended to the case of the spectrum (1) and the expression of the free energy F takes the form

$$F=-T\sum_{\mathbf{q},q_\perp}\ln\frac{\pi T}{\alpha\left[\epsilon+\eta_2\left(\mathbf{q}^2+\frac{4J^2}{v_F^2}\sin^2(q_\perp s/2)\right)\right]}, \qquad (2)$$

where α and

$$\eta_2=-\frac{v_F^2\tau^2}{2}\left[\psi\left(\frac{1}{2}+\frac{1}{4\pi\tau T}\right)-\psi\left(\frac{1}{2}\right)-\frac{1}{4\pi\tau T}\psi'\left(\frac{1}{2}\right)\right] \qquad (3)$$

are the coefficients of two-dimensional GL theory, $\psi(x)$ is the digamma function, and τ is the electron relaxation time. It is essential to note that $\epsilon=\ln(T/T_c)\simeq(T-T_c)/T_c$ is supposed to be small ($\epsilon\ll 1$) in the expression (2) and this hypothesis provides the convergency of the summation over in-plane momenta in the region of small $q\ll\sqrt{T/D_\parallel}$.

The crossover temperature is defined by the analysis of the denominator in the expression (2). Evidently, in the strictly two-dimensional case, the hopping integral $J\rightarrow 0$ and the integration over $q_\perp$ is reduced to the redefinition of the density of states only. The Lawrence-Doniach crossover in the fluctuation behavior takes place when the integration over the transverse variable becomes important and changes the temperature dependence of the appropriate fluctuation correction. It is easy to see that this happens when the transversal part reaches the reduced temperature ϵ [or, in other words, when $\xi_c(\epsilon_{LD})\approx s$]:

$$\epsilon_{LD}=\frac{4\eta_2J^2}{v_F^2}. \qquad (4)$$

The observation of another crossover is the essence of this paper and it is related to the breakdown of the GL approach in the description of fluctuations relatively far from the transition: in this case ϵ is not a small parameter, so the short-wavelength and the high-frequency fluctuations have to be taken into account. Formally speaking, the ''fluctuation propagator,'' which in the GL approach is proportional to the simplified expression $[\epsilon+\eta_2\mathbf{q}^2+\epsilon_{LD}\sin^2(q_\perp s/2)]^{-1}$ and is valid for the long-wavelength approximation only, has to be substituted in Eq. (2) by its more general form:[6]

$$L(\mathbf{q},q_\perp,\omega_\mu)=-\rho\left[\epsilon+\psi\left(\frac{1}{2}+\frac{\omega_\mu}{4\pi T}+\frac{2\eta_2}{\pi^2}\times\left(\mathbf{q}^2+\frac{4J^2}{v_F^2}\sin^2\frac{q_\perp s}{2}\right)\right)-\psi\left(\frac{1}{2}\right)\right]^{-1}, \qquad (5)$$

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and the summation over all bosonic Matsubara frequencies has to be accomplished side by side with the integration over all momenta [$\rho = m/(2\pi s)$ is the single-spin quasiparticle normal density of states].

The appropriate generalization of the Aslamazov-Larkin theory of paraconductivity[7] for the high-temperature region was carried out in Ref. 8 where was found the formula for paraconductivity expressed in terms of the universal function $f(\epsilon)$:

$$\sigma_{fl} = \frac{e^2}{16\hbar s} f(\epsilon). \qquad (6)$$

In the GL region of temperature, where $\epsilon \ll 1$, $f(\epsilon) = 1/\epsilon$ and the result naturally coincides with the well-known Aslamazov-Larkin one. In the opposite case $\epsilon \gg 1$, for clean two-dimensional superconductors it was found that $f(\epsilon) \sim 1/\epsilon^3 = 1/\ln^3(T/T_c)$. Let us stress that, even in its generalized form, the two-dimensional paraconductivity is universal; it contains the only intrinsic parameter s as prefactor and this fact permits us to rescale the results and to compare them for different compounds.

In the theoretical consideration it was natural to assume formally the very rigid restriction $\epsilon \gg 1$ for the validity of the high-temperature asymptotic behavior. Nevertheless, as will be seen below, in experiments the crossover to this asymptotic behavior takes place *universally* for all the investigated samples at $\epsilon \sim 0.23$ and this can be attributed to some particularly fast convergence of the integrals in the expression of $f(\epsilon)$.

The long tails in the in-plane fluctuation conductivity of HTS materials have been observed frequently. One of the efforts to fit the high-temperature paraconductivity with the extended Aslamazov-Larkin (AL) theory results was undertaken in Ref. 9 where the deviation of the excess conductivity from AL behavior was analyzed for three $Bi_2Sr_2CaCu_2O_8$ epitaxial films. A very good fit with the formula (6) was found in the region of temperatures $0.02 \lesssim \epsilon \lesssim 0.14$. We show here that the careful analysis of the higher temperature region (just above the edge of the region investigated in Ref. 9) allows us to observe the surprisingly early approaching to the short-wavelength fluctuation (SWF) asymptotic regime [at the reduced temperature $\epsilon^* \sim \ln(T^*/T_c) \sim 0.23$].

We have performed resistivity measurements of three different HTS compounds: a melt-textured $YBa_2Cu_3O_7$ sample (Y123), a $Bi_2Sr_2CaCu_2O_8$ (Bi2212) thick film, and a highly textured $Bi_2Sr_2Ca_2Cu_3O_{10}$ (Bi2223) tape. The Y123 was obtained by melting;[10] the sample was cut in a nearly regular parallelepipedal shape with a cross section of about 4 mm^2 and a length of 7 mm. The resistivity measurements were performed from 85 to 330 K and are shown in Fig. 1(a). The critical temperature, defined as the point where the temperature derivative is maximal, is 92 K; $\rho_N(100\ K) = 180\ \mu\Omega$ cm, where ρ_N is the resistivity value extrapolated from the high-temperature region where the normal state ρ is linear. The Bi2212 film was prepared by a liquid phase epitaxy technique.[11] The film has a thickness of about 1 μm. The resistivity measurements were performed from 80 to 170 K [see Fig. 1(b)]. The critical temperature was estimated to be 84.2 K and $\rho_N(100\ K) = 150\ \mu\Omega$ cm. The Bi2223 tape was obtained by means of the power in tube procedure, as described elsewhere.[12] The thickness of the oxide filament inside the tape was about 30 μm; the filament turned out to be strongly textured (rocking angle $\approx 8°$) with the c axis oriented perpendicular to the tape plane. The resistivity measurements were performed in the range from 100 to 250 K, after removing the silver sheathing chemically, and are reported in Fig. 1(c). The critical temperature was estimated to be 108 K and $\rho_N(100\ K) = 300\ \mu\Omega$ cm. We ascribe this high value of ρ_N to different causes: first, the grain boundaries may determine a resistance in series with the grain resistance; second, the chemical treatment may have damaged the surface of the sample and the effective cross section of the superconductor can be decreased.

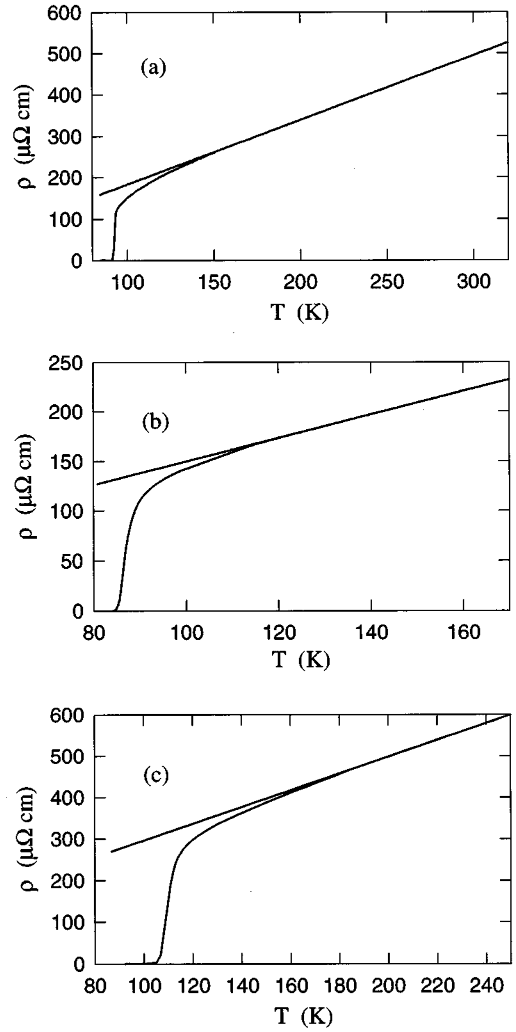

FIG. 1. ρ versus T for (a) Y123, (b) Bi2212, (c) Bi2223. In the figures ρ_N (obtained best fitting the data in the ranges indicated in the text) are also shown.

The excess conductivity was estimated by subtracting the background of the normal-state conductivity $\sigma_N = 1/\rho_N$. The evaluation of ρ_N was made with particular accuracy; in fact, starting the interpolation at a certain temperature corresponds to forcing σ_{fl} to vanish artificially at such temperature. Therefore, we need to estimate ρ_N at a temperature as large as possible and to verify that ρ_N does not depend on the temperature range where the interpolation is performed. In the case of the Y123 sample the resistivity shows a linear

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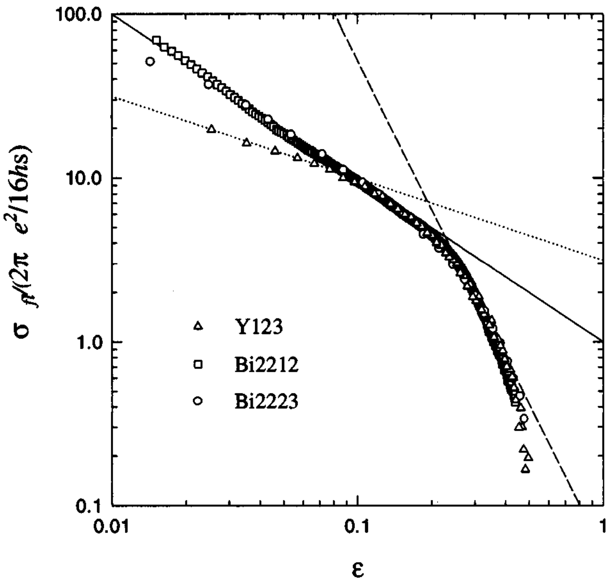

FIG. 2. $\sigma_{fl}(16hs/2\pi e^2)$ vs ϵ for Y123 (triangle), Bi2212 (square), and Bi223 (circlet); the solid line is $1/\epsilon$, the dashed line $0.055/\epsilon^3$, and the dotted line is $3.2/\sqrt{\epsilon}$.

behavior from 160 to 330 K. In this range we have verified that ρ_N does not change by shifting the interpolation temperature region. The ρ_N line obtained by a best fit in the temperature range 160 to 330 K is also reported in Fig. 1(a). Therefore, for the Y123 sample, the upper limit of ϵ at which the excess conductivity may be analyzed is $\epsilon_{up} \approx \ln(160/92) = 0.55$. In an analogous way we obtain $\epsilon_{up} \approx 0.46$ and 0.51 for Bi2212 and Bi2223, respectively. The ρ_N lines obtained by best fitting the high-temperature resistivity data are also reported for Bi2212 and Bi2223 in Figs. 1(b) and 1(c), respectively. In Fig. 2, in a log-log scale, we plot $\sigma_{fl}(16hs/2\pi e^2)$ as a function of ϵ for the three samples; the solid line represents $1/\epsilon$, the dashed line $0.055/\epsilon^3$, and the dotted line is $3.2/\sqrt{\epsilon}$. The interlayer distance s is considered as a free parameter and it has been adjusted so that the experimental data can follow the $1/\epsilon$ behavior in the ϵ region where the AL behavior is expected.

We can see that all the curves exhibit the same general behavior. The region where the two-dimensional $1/\epsilon$ behavior is followed has different extensions for each compound, depending on its anisotropy, and at $\epsilon \approx 0.23$ all the curves bend downward and follow the same asymptotic $1/\epsilon^3$ behavior.

We discuss now some features in detail: (1) The interlayer distance values we find are the following: for Y123 we obtain $s \approx 13$ Å which must be compared with the YBCO interlayer distance that is about 12 Å; for Bi2212 we obtain $s \approx 11$ Å to be compared with 15 Å; and for Bi2223 we obtain $s \approx 25$ Å to be compared with 18 Å. The differences in the interlayer distance evaluation are all compatible with the uncertainty on the geometrical factors. We point out that the smallest error is for Y123 (about 10%) that is a bulk sample with a well-defined geometry. Larger errors are found for the Bi2212 thick film (about 30%), for which the evaluation of the thickness is rough, and for the Bi2223 tape (about 40%), for which an overestimation of the cross section of the tape is possible, as we mentioned above. We conclude that the AL behavior is well followed. (2) On the low ϵ value side ($\epsilon < 0.2$) the three compounds show different behaviors due to the different extension of the AL region. The least anisotropic compound, Y123, for $\epsilon < 0.1$ bends going asymptotically to the 3D behavior ($1/\epsilon^{0.5}$) showing the LD crossover at $\epsilon \approx 0.09$; the Bi2223 sample starts to bend for $\epsilon < 0.03$ while the most anisotropic Bi2212 in the overall ϵ range considered shows the two-dimensional behavior. (3) On the high ϵ value side, starting from the AL behavior, the curves show a crossover at about $\epsilon = 0.23$ and then bend downward following the asymptotic $1/\epsilon^3$ behavior. At the value $\epsilon \approx 0.45$ all the curves drop indicating the end of the observable fluctuation regime. This value is lower than the above reported ϵ_{up} values, at which the fluctuation conductivity comes out to be zero.

To conclude we have observed in three different HTS compounds the universal high-temperature behavior of the in-plane conductivity that manifests itself in the two-dimensional regime, once the LD crossover is passed. Beyond the AL regime all the curves soon reach the SWF $1/\epsilon^3$ regime. For all the compounds the crossover occurs at the same point $\epsilon \approx 0.23$, which corresponds to $T \approx 1.3T_c$ and, therefore, is experimentally well observable. The universality of the paraconductivity behavior is much more surprising if we consider that it has been observed in three compounds with different crystallographic structures and anisotropies, and moreover prepared by means of very different techniques.

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