Universal scaling behavior of the c-axis resistivity of high-temperature superconductors

Y. H. Su
Center for Advanced Studies, Tsinghua University, Beijing 100084, China

H. G. Luo
Institute of Theoretical Physics and Interdisciplinary Centre of Theoretical Studies, Chinese Academy of Sciences, P.O. Box 2735, Beijing 100080, China

T. Xiang
Institute of Theoretical Physics and Interdisciplinary Centre of Theoretical Studies, Chinese Academy of Sciences, P.O. Box 2735, Beijing 100080, China

We propose and show that the c-axis transport in high-temperature superconductors is controlled by the pseudogap energy and the c-axis resistivity satisfies a universal scaling law in the pseudogap phase. We derived approximately a scaling function for the c-axis resistivity and found that it fits well with the experimental data of Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$, Bi$_2$Sr$_2$CaCu$_2$O$_{10+\delta}$, and YBa$_2$Cu$_3$O$_{7-\delta}$. Our works reveals the physical origin of the semiconductor-like behavior of the c-axis resistivity and suggests that the c-axis hopping is predominantly coherent.

I. INTRODUCTION

In the pseudogap phase of high-temperature (high-$T_c$) superconductors, the temperature dependence of the c-axis resistivity $\rho_c$ is semiconductor-like ($d\rho_c/dT < 0$), in contrast to the in-plane resistivity $\rho_{ab}$ whose temperature dependence is metal-like ($d\rho_{ab}/dT > 0$). This dramatic difference between $\rho_c$ and $\rho_{ab}$ is not what one might expect within the conventional Fermi liquid theory. It has stimulated vast theoretical and experimental investigations on the interlayer dynamics of high-$T_c$ cuprates. However, the existing theoretical models based on the notion of dynamic confinement or incoherent interlayer hopping could not give a natural and unified explanation for experimental data.

In high-$T_c$ cuprates, the c-axis hopping integral is highly anisotropic and depends strongly on the in-plane momenta. Based on electron structure calculations and symmetry analysis, it was shown that the c-axis hopping integral in high-$T_c$ cuprates with tetragonal symmetry is given by

$$t_c(k) \propto (\cos k_x - \cos k_y)^2, \quad (1)$$

This peculiar in-plane momentum dependence of $t_c(k)$ results from the hybridization between the bonding O 2$p$ orbitals and Cu 4$s$ or $3d_{3z^2-r^2}$ orbitals in each CuO$_2$ plane. It was confirmed by the angle-resolved photoemission spectroscopy (ARPES) measurements. Along the two diagonals of the Brillouin zone, these are also the directions where the nodes of the $d_{z^2}$-wave superconducting gap or the normal state pseudogap (it has also the $d$-wave symmetry) are located. Thus the quasiparticle excitations around the gap nodes have almost no contributions to the c-axis hopping. This indicates that the c-axis dynamics is governed by the quasiparticle excitations around the antinodal points.

The interplay between $t_c(k)$ and the $d$-wave energy gap can affect strongly the c-axis dynamics of electrons. In the superconducting state, it leads to, for example, a $T^5$ temperature dependent c-axis superfluid density in low temperatures, in contrast to the linear-$T$ behavior of the in-plane superfluid density. In a pseudogap normal state, since the contribution from the quasiparticle excitations around the nodal points is suppressed by the integral (1), the interlayer transport would behave as in a gapped system and $\rho_c$ should show a thermally activated behavior. On the contrary, the in-plane dynamics is governed by low-lying excitations around the gap nodes in low temperatures. Thus $\rho_{ab}$ would behave similarly as in a gapless system. This explains naturally why the temperature dependence of $\rho_c$ is semiconductor-like while $\rho_{ab}$ is metal-like in the pseudogap phase. In making this argument, we have implicitly assumed that the c-axis hopping is predominantly coherent. We believe this assumption is correct, at least in the limit of weak impurity scattering. It is supported by the observation of the bilayer splitting of the bonding and antibonding bands in the ARPES spectra, the intrinsic $T^5$ temperature dependence of the c-axis superfluid density, as well as the recent angular-dependent magnetoresistance measurement.

The c-axis dynamics could be also affected by phonon or other collective excitations. Along the nodal directions, a dispersion kink has been observed in the spectral function of electrons by the ARPES in both superconducting and normal states. This kink is believed to result from the electron-phonon coupling. Along the antinodal direction, a stronger kink feature was observed in the single-electron spectrum below $T_c$. This stronger kink seems to be correlated with the peak-dip hump lineshape as observed in the ARPES spectra at antinodal points. It is likely to be due to the coupling of electrons with a magnetic resonance mode. It could...
also be explained by the coupling of electrons with a B1g-phonon mode. However, this coupling kink together with the peak-dip-lump structure in the ARPES spectra disappears above Tc. It suggests that the coupling between electrons and phonons or magnetic resonance modes is strong around the gap nodes but very weak along the antinodal directions in the normal state. Thus we believe that the coupling between electrons and phonons or magnetic resonance modes should have very weak or negligible effect on c-axis transport properties in the normal state.

II. SCALING HYPOTHESIS

The above discussion indicates that in the normal state the pseudogap is the only energy scale governing the c-axis hopping around the antinodal regions. If the c-axis hopping is predominately coherent, this would suggest this axis hopping around the antinodal regions. If the c-axis the pseudogap is the only energy scale governing the c-axis transport properties in the normal state.

Thus we believe that the coupling between electrons and phonons or magnetic resonance modes is strong around the gap nodes but very weak along the antinodal directions in the normal state. It can be obtained by evaluating the current-current correlation function in the linear response theory. In the highly anisotropic cuprate superconductors, the c-axis hopping can be considered as perturbation. In this case, the c-axis conductivity \( \sigma_c \) is given by

\[
\sigma_c(T) \propto -\alpha_c g \left( \frac{T}{\Delta} \right),
\]

where \( g(x) \) is a scaling function, \( \alpha_c \) is a doping-dependent coefficient, and \( \Delta \) is the maximal value of the pseudogap. The detailed form of \( g(x) \) is unknown. However, as will be shown below, an approximate but accurate expression for \( g(x) \) can be obtained by simple theoretical analysis.

\( \rho_c \) is determined by the interlayer hopping integral and the scattering mechanism. It can be obtained by evaluating the current-current correlation function in the linear response theory. In the highly anisotropic cuprate superconductors, the c-axis hopping can be considered as perturbation. In this case, the c-axis conductivity \( \sigma_c \) is given by

\[
\sigma_c(T) \propto -\sum_k \int d\omega v_c^2(k) A^2(k,\omega) \frac{\partial f(\omega)}{\partial \omega},
\]

where \( v_c(k) \propto t_c(k) \) is the c-axis velocity of electrons and \( f(\omega) \) is the Fermi function. \( A(k,\omega) \) is the spectral function of pseudogapped electrons. At present, an accurate microscopic description of \( A(k,\omega) \) is not available since the mechanism of pseudogap is still unclear. However, phenomenologically one can assume \( A(k,\omega) \) to have the form

\[
A(k,\omega) = \frac{1}{\pi} \frac{\Gamma}{(\omega - E_k)^2 + \Gamma^2},
\]

where \( E_k = \pm \sqrt{\varepsilon_k^2 + \Delta^2 \cos^2(2\phi_k)} \) is the energy dispersion of pseudogapped electrons, \( \varepsilon_k = \arctan(k_x/k_y) \) and \( \varepsilon_k \) is the usual normal state energy measured from the Fermi level, \( \Gamma \) is the linewidth, proportional to the scattering rate. In the normal state above the pseudogap phase, \( \rho_{ab} \) varies linearly with T. It suggests that \( \Gamma \) scales linearly with T. This linear T dependent \( \Gamma \) is a basic assumption made in the marginal Fermi liquid theory. It gives a natural account of many experimental results, including the linear in-plane resistivity. Here we will also adopt this assumption. Substituting Eq. \( 11 \) into Eq. \( 3 \), it can be shown that to the leading order approximation in \( \Gamma \), \( \sigma_c \) is given by

\[
\sigma_c(T) \propto -\frac{1}{T} \int d\omega N_c(\omega) \frac{\partial f(\omega)}{\partial \omega},
\]

where \( N_c(\omega) \) is an effective measure of the c-axis tunneling probability of electrons

\[
N_c(\omega) = \frac{1}{V} \sum_k t_c^2(k) \delta(\omega - E_k),
\]

and \( V \) is the system volume. \( N_c(\omega) \) reduces to the density of states \( N(\omega) \) of d-wave gapped quasiparticles if \( t_c(k) \) in Eq. \( 3 \) is set to 1. Figure \ref{fig:1} compares the energy dependence of \( N_c(\omega) \) with \( N(\omega) \). It shows clearly that the c-axis hopping of low-energy quasiparticles is strongly suppressed by \( t_c(k) \) and the low energy \( N_c(\omega) \) is vanishingly small compared with that at \( \omega \sim \Delta \). Thus the temperature considered is not too low compared with \( \Delta \), the contribution of low energy electrons to the c-axis hopping can be ignored. This has motivated us to replace approximately \( N_c(\omega) \) by a step function,

\[
N_c(\omega) \sim N_0 \theta(\omega - \Delta),
\]

where \( N_0 \) is an average value of \( N_c(\omega) \) at \( \omega \gtrsim \Delta \). Here \( \Delta \) should be slightly smaller than the true maximum pseudogap since the contribution from the states with \( \omega < \Delta \) should be effectively included in the above approximation. Apparently this is a crude approximation, but it allows us to get an analytical expression for \( \sigma_c \). By substituting \ref{eq:7} into \ref{eq:5}, we find that \( \rho_c \) is approximately given by

\[
\rho_c(T) = \sigma_c^{-1}(T) \approx \frac{\alpha_c T}{\Delta} \exp \left( \frac{\Delta}{T} \right)
\]
in the limit $T_c \ll T \ll \Delta$. We have tested this formula with numerical calculations without taking the step-function approximation for $N_c(\omega)$. We find that Eq. (8) does fit the numerical curves very well in the above limit. Though this formula is derived within a limited temperature regime, it captures qualitatively the main features of $\rho_c$ in the whole temperature range in the normal state. In the low temperature limit $T \ll \Delta$, Eq. (8) is thermally activated. In the high temperature limit $T \gg \Delta$, $\rho_c$ varies linearly with $T$. Both agree with the experimental observations.

III. COMPARISON WITH EXPERIMENTS

To test our single-parameter scaling conjecture, we have analyzed the experimental data of $\rho_c$ published for Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ (Bi2212) by Watanabe et al.\textsuperscript{19,20} by Chen et al.\textsuperscript{21} and by Guira et al.\textsuperscript{22} for Bi$_2$Sr$_2$Ca$_2$Cu$_3$O$_{10+\delta}$ (Bi2223) by Fujii et al.\textsuperscript{23} and for YBa$_2$Cu$_3$O$_{7-\delta}$ (Y123) by Yan et al.\textsuperscript{24} and by Babic et al.\textsuperscript{25} We find that all these experimental data of $\rho_c$ can be well fitted by Eq. (8) nearly in the entire temperature range. Fig. 2(a)-(c) show the scaling behavior of $\rho_c$ for Bi2212, Bi2223, and Y123, respectively. For Bi2212, we only showed the data published by Watanabe et al.\textsuperscript{26} for clarity. The experimental data of Bi2212 by the same group\textsuperscript{26} can be also scaled onto the same curve. The experimental data of Bi2212 published by other two groups\textsuperscript{21,22} are showed in Fig. 3. The fitting parameters $\Delta$ and $\alpha_c$ for Bi2212, Bi2223, and Y123 at different doping levels are given in Table I-V. The doping-dependence of $\Delta$ and comparison with the ARPES data will be discussed later.

![Graphs showing scaling behaviors of $\rho_c$ and $[\rho_{ab}(T) - \rho_{ab}(0)]/(\alpha_{ab}T)$ for Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$, Bi$_2$Sr$_2$Ca$_2$Cu$_3$O$_{10+\delta}$, and YBa$_2$Cu$_3$O$_{7-\delta}$](image)

**FIG. 2:** Scaling behaviors of the c-axis and the in-plane resistivities. (a)-(c), the scaling curves of $\rho_c/\alpha_c$ versus $T/\Delta$ for the experimental data published by Watanabe et al.\textsuperscript{26} for Bi2212, by Fujii et al.\textsuperscript{23} for Bi2223, and by Yan et al.\textsuperscript{24} ($\delta = 0.3$) and by Babic et al.\textsuperscript{25} for Y123, respectively. The parameters $\alpha_c$ and $\Delta$ are determined by fitting the experimental data with Eq. (8) well above $T_c$. The dashed lines in (a)-(c) denote the scaling function $g(x) = x \exp(1/x)$. (d)-(f), the corresponding normalized in-plane resistivities $[\rho_{ab}(T) - \rho_{ab}(0)]/(\alpha_{ab}T)$ as a function of $T/\Delta$ is shown for Bi2212, Bi2223, and Y123, respectively. Different curves are shifted vertically from each other for clarity. Both $\alpha_{ab}$ and $\rho_{ab}(0)$ can be determined by fitting the experimental data in the linear regime of $\rho_{ab}$ with Eq. (10). In this linear regime, the ratio $[\rho_{ab}(T) - \rho_{ab}(0)]/(\alpha_{ab}T)$ is equal to 1 within measurement errors. However, below a doping dependent temperature $T^*$, which is commonly defined as the onset temperature of the pseudogap, $\rho_{ab}$ begins to deviate from this linear-T behavior and $[\rho_{ab}(T) - \rho_{ab}(0)]/(\alpha_{ab}T)$ drops quickly with decreasing temperature, as marked roughly by the arrows in (d)-(f).
The measurement data deviate slightly from the scaling curve near $T_c$. This is due to the superconducting fluctuations. This deviation is not unexpected since in this case the single-parameter scaling law (2) should be modified to include the contributions from the superconducting fluctuations. The above analysis shows that the scaling behavior of $\rho_c$ is universal and the scaling function $g(x)$ is approximately given by

$$g(x) = x \exp \left( \frac{1}{x} \right)$$  \hspace{1cm} (9)$$

independent of the doping concentration as well as the chemical structures for these multilayer compounds. Thus there is only one energy scale characterizing the $c$-axis dynamics in the pseudogap phase. It suggests that the pseudogap is the only energy scale governing the quasiparticle excitations at the antinodal points.

**TABLE I:** The fitting parameters $\Delta$, $\alpha_c$, $\alpha_{ab}$, and $\rho_{ab}(0)$ for Bi2212 published in Ref. [21] (the first 7 rows) and Ref. [24] (the rest 4 rows). The units of $\alpha_{ab}$ and $\rho_{ab}(0)$ are $10^{-6}$Ωcm/K and $10^{-4}$Ωcm, respectively.

| $\delta$ | $T_c$ (K) | $\Delta$ (K) | $\alpha_c$(Ωcm) | $\alpha_{ab}$ | $\rho_{ab}(0)$ |
|----------|-----------|--------------|------------------|--------------|----------------|
| 0.2135   | 71        | 416.6        | 2.74             | 2.60         | 1.3            |
| 0.216    | 77        | 397.4        | 2.56             | -            | -              |
| 0.22     | 83        | 367.3        | 2.31             | 1.78         | 1.1            |
| 0.23     | 86        | 344.5        | 1.95             | -            | -              |
| 0.26     | 88        | 232.9        | 1.07             | 1.03         | 0.22           |
| 0.27     | 85        | 207.1        | 0.82             | 0.78         | 0.18           |
| 0.28     | 79        | 143.3        | 0.56             | -            | -              |
| 0.25     | 90        | 255          | 1.14             | -            | -              |
| 0.26     | 86        | 217          | 0.91             | -            | -              |
| 0.27     | 83        | 184.7        | 0.67             | -            | -              |
| 0.28     | 79        | 150.1        | 0.34             | -            | -              |

It is interesting to compare the above result with the temperature dependence of $\rho_{ab}$. In high temperatures, $\rho_{ab}$ shows linear temperature dependence

$$\rho_{ab}(T) = \alpha_{ab}T + \rho_{ab}(0), \hspace{1cm} (10)$$

where $\alpha_{ab}$ is the slope of the linear resistivity and $\rho_{ab}(0)$ is the zero temperature resistivity extrapolated from high temperature data of $\rho_{ab}$. Fig. 2(d)-(f) show the experimental data of $[\rho_{ab}(T) - \rho_{ab}(0)]/(\alpha_{ab}T)$ as a function of $T/\Delta$ for Bi2212, Bi2223, and Y123, respectively. The slope $\alpha_{ab}$ and $\rho_{ab}(0)$ for Bi2212, Bi2223, and Y123 are given in Table I, IV, and V, respectively. Here $\Delta$ is the pseudogap value determined from $\rho_c$. An intriguing result revealed by Fig. 2(d)-(f) is that $T^*$ (arrows in the figure) is approximately proportional to $\Delta$, i.e., $T^* \propto \Delta$. This result is reminiscent of the relationship between the superconducting transition temperature $T_c$ and the superconducting energy gap $\Delta_0$. For a BCS mean-field d-wave superconductor, $T_c/\Delta_0 \sim 0.47$. Here the ratio $T^*/\Delta$ is about 0.55 for Bi2212, 0.66 for Bi2223, and 0.71
for Y123. These values are larger than the BCS value. This is not unexpected since, as pointed out above, \( \Delta \) obtained by Eq. 5 is smaller than the true pseudogap.

Figure 4 shows the doping dependence of \( \Delta \) and \( T_c^{\star} \) for Bi2212. For comparison, the ARPES data published by Campuzano et al. [22] for Bi2212 are also shown in the figure. The data of \( T_c^{\star} \) obtained agree well with the ARPES data. Our results of \( \Delta \) give a lower bound for the maximum pseudogap as expected. They are less fluctuating than the ARPES data. It suggests that \( \rho_c \) is a good probe for the pseudogap.

The above comparison shows that \( \rho_c \) satisfies a simple single-parameter scaling law. We believe that the scaling function \( g(x) \) is universal and given approximately by Eq. 19 for all high-\( T_c \) cuprates whose dominant c-axis hopping integral between two neighboring CuO\(_2\) layers is coherent and has the in-plane momentum dependence given by Eq. 11. This includes all multilayer compounds and some single-layer compounds in which Cu atoms in the two neighboring unit cells lie collinearly along the c-axis, such as HgBa\(_2\)CuO\(_{4-\delta}\). However, for other single-layer compounds, for example Bi\(_2\)Sr\(_2\)CuO\(_{6+\delta}\) and La\(_{2-\delta}\)Sr\(_{\delta}\)CuO\(_4\), Cu atoms of two adjacent CuO\(_2\) planes do not lie collinearly along the c-axis. In this case, the c-axis hopping integral has the form [25]:

\[
t_c(k) \propto \frac{k_x}{2} \cos \frac{k_y}{2} (\cos k_x - \cos k_y)^2.
\]

### Table III: The fitting parameters \( \Delta \) and \( \alpha_c \) for Bi2212 published by Giura et al. [22].

| \( \delta \) | \( T_c (K) \) | \( \Delta (K) \) | \( \alpha_c (\Omega cm) \) |
|---|---|---|---|
| 0.232 | 87.4 | 321.89 | 1.99 |
| 0.240 | 88.1 | 281.62 | 1.64 |
| 0.253 | 92.0 | 263.84 | 1.17 |
| 0.260 | 87.0 | 194.8 | 0.69 |
| 0.265 | 85.9 | 192.96 | 0.54 |
| 0.267 | 86.2 | 191.69 | 0.34 |
| 0.271 | 87.6 | 177.88 | 0.30 |

### Table IV: The fitting parameters \( \Delta, \alpha_c, \alpha_{ab}, \) and \( \rho_{ab}(0) \) for Bi2223 published in Ref. [22]. The units of \( \alpha_{ab} \) and \( \rho_{ab}(0) \) are \( 10^{-6} \Omega cm/K \) and \( 10^{-5} \Omega cm \), respectively.

| Samples | \( T_c (K) \) | \( \Delta (K) \) | \( \alpha_{c}(\Omega cm) \) | \( \alpha_{ab} \) | \( \rho_{ab}(0) \) |
|---|---|---|---|---|---|
| a | 87 | 403.29 | 7.52 | - | - |
| b | 93 | 378.74 | 6.18 | 2.34 | -2.21 |
| c | 98 | 361.74 | 5.38 | - | - |
| d | 104 | 330.17 | 4.15 | 2.21 | -4.45 |
| e | 106 | 300.65 | 3.4 | 2.06 | -5.24 |
| f | 108 | 272.72 | 2.78 | 1.85 | -5.81 |
| g | 108 | 245.46 | 2.28 | 1.77 | -6.02 |
| h | 108 | 221.42 | 1.85 | 1.57 | -4.77 |
| i | 108 | 212.85 | 1.03 | - | - |

### Table V: The fitting parameters \( \Delta, \alpha_c, \alpha_{ab}, \) and \( \rho_{ab}(0) \) for Y123 published in Ref. [22] (\( \delta = 0.3 \)) and Ref. [25] (rest).

| \( \delta \) | \( T_c (K) \) | \( \Delta (K) \) | \( \alpha_c(m\Omega cm) \) | \( \alpha_{ab}(\mu\Omega cm/K) \) | \( \rho_{ab}(0)(\mu\Omega cm) \) |
|---|---|---|---|---|---|
| 0.35 | 61.4 | 317.49 | 10.41 | 1.36 | 1.45 |
| 0.21 | 83.7 | 221.87 | 3.83 | 0.86 | 4.77 |
| 0.15 | 91.0 | 194.25 | 2.48 | 0.84 | 4.84 |
| 0.30 | 65.0 | 318.21 | 105.01 | - | - |

### IV. SUMMARY

The interplay between the anisotropic c-axis hopping integral and the pseudogap has profound consequence on the c-axis dynamics. It leads to a universal scaling behavior for \( \rho_c \) in the pseudogap phase. The excellent agree-
ment between our theoretical analysis and experimental data suggests that the pseudogap is the only energy scale governing the quasiparticle excitations around the antinodal points in the normal state of high-\(T_c\) materials.

Our analysis for the scaling behavior of \(\rho_c\) is valid generally, independent of the mechanism of normal state pseudogap. We believe that it can be generalized and applied to other c-axis transport quantities. Before closing the paper, we would like to make a general conjecture: if \(F_{c,e}\) is the electronic contribution to a c-axis transport coefficient, then \(F_{c,e}\) should satisfy a single-parameter scaling law

\[
F_{c,e}(T) = \beta_c f_c \left( \frac{T}{\Delta} \right)
\]  

at least at not too low temperatures in the pseudogap phase. Here \(\beta_c\) is a doping-dependent coefficient and \(f_c(x)\) is a corresponding scaling function. A thorough examination of this conjecture, no matter being approved or disapproved, would shed light on further understanding of properties of quasiparticle excitations around the antinodal points. This will help us to understand more about the mechanism of the high-\(T_c\) superconductivity since the superconducting pairing is strongest at the antinodal points.

Acknowledgments

We are grateful to T. Watanabe, X. H. Chen, and S. Sarti for kindly sending us their experimental data published in Refs. [19, 20, 21, 22, 23] respectively. We thank L. Yu for valuable discussions. The work was supported by the National Natural Sciences Foundation of China.

1. P. W. Anderson, *The theory of superconductivity in the high-\(T_c\) cuprates* (Princeton University Press, Princeton, New Jersey, 1997).
2. N. Kumar and A. M. Jayannavar, Phys. Rev. B *45*, 5001 (1992).
3. M. Turlakov and A. J. Leggett, Phys. Rev. B *63*, 064518 (2001).
4. O. K. Andersen, A. I. Liechtenstein, O. Jepsen, and F. Paulsen, J. Phys. Chem. Solids *56*, 1573 (1995).
5. S. Chakravarty, A. Sudbø, P. W. Anderson, and S. Strong, Science *261*, 337 (1993).
6. T. Xiang and J. M. Wheatley, Phys. Rev. Lett. *77*, 4632 (1996).
7. T. Xiang, C. Panagopoulos, and J. R. Cooper, Int. J. Mod. Phys. B *12*, 1007 (1998).
8. D. L. Feng, N. P. Armitage, D. H. Lu, A. Damascelli, J. P. Hu, P. Bogdanov, A. Lanzara, F. Ronning, K. M. Shen, H. Eisaki, C. Kim, J.-i. Shimoymama, K. Kishio, and Z.-X. Shen, Phys. Rev. Lett. *86*, 5550 (2001).
9. A. G. Loeser, Z.-X. Shen, D. S. Dessau, D. S. Marshall, C. H. Park, P. Fournier, and A. Kapitulnik, Science *273*, 325 (1996).
10. H. Ding, T. Yokoya, J. C. Campuzano, T. Takahashi, M. Randeria, M. R. Norman, T. Mochiku, K. Kadowaki, and J. Giapintzakis, Nature *382*, 51 (1996).
11. A. S. Alexandrov, V. V. Kabanov, and N. F. Mott, Phys. Rev. Lett. *77*, 4796 (1996).
12. N. E. Hussey, M. Abdel-Jawad, A. Carrington, A. P. Mackenzie, and L. Balicas, Nature *425*, 814 (2003).
13. P. V. Bogdanov, A. Lanzara, S. A. Kellar, X. J. Zhou, E. D. Lu, W. J. Zheng, G. Gu, J.-i. Shimoymama, K. Kishio, H. Ikeda, R. Yoshizaki, Z. Hussain, and Z.-X. Shen, Phys. Rev. Lett. *85*, 2581 (2000).
14. A. Lanzara, P. V. Bogdanov, X. J. Zhou, S. A. Kellar, D. L. Feng, E. D. Lu, T. Yoshida, H. Eisaki, A. Fujimori, K. Kishio, J.-i. Shimoymama, T. Noda, S. Uchida, Z. Hussain, and Z.-X. Shen, Nature *412*, 510 (2001).
15. A. D. Gromko, A. V. Fedorov, Y.-D. Chuang, J. D. Koralek, Y. Aiura, Y. Yamaguchi, K. Oka, Yoichi Ando, and D. S. Dessau, Phys. Rev. B *68*, 174520 (2003).
16. J. C. Campuzano, H. Ding, M. R. Norman, H. M. Fretwell, M. Randeria, A. Kaminski, J. Mesot, T. Takeuchi, T. Sato, T. Yokoya, T. Takahashi, T. Mochiku, K. Kadowaki, P. Gupatasarima, D. G. Hinks, Z. Konstantinovic, Z. Z. Li, and H. Raffy, Phys. Rev. Lett. *83*, 3709 (1999).
17. T. Cuk, F. Baumberger, D. H. Lu, N. Ingle, X. J. Zhou, H. Eisaki, N. Kaneko, Z. Hussain, T. P. Devereaux, N. Nagosa, and Z.-X. Shen, Phys. Rev. Lett. *93*, 117003 (2004).
18. C. M. Varma, P. B. Littlewood, S. Schmitt-Rink, E. Abraham, and A. E. Ruckenstein Phys. Rev. Lett. *63*, 1996 (1989).
19. T. Watanabe, T. Fujii, and A. Matsuda, Phys. Rev. Lett. *79*, 2113 (1997); Recent Res. Devel. Physics *5*, 51 (2004).
20. T. Watanabe, T. Fujii, and A. Matsuda, Phys. Rev. Lett. *84*, 5548 (2000).
21. X. H. Chen, M. Yu, K. Q. Ruan, S. Y. Li, Z. Gui, G. C. Zhang, and L. Z. Cao, Phys. Rev. B *58*, 14219 (1998).
22. M. Giura, R. Fastampa, S. Sarti, and E. Silva, Phys. Rev. B *68*, 134505 (2003).
23. T. Fujii, I. Terasaki, T. Watanabe, and A. Matsuda, Phys. Rev. B *66*, 024507 (2002).
24. Y. F. Yan, P. Matl, J. M. Harris, and N. P. Ong, Phys. Rev. B *52*, R751 (1995).
25. D. Babic, J. R. Cooper, J. W. Hodby, and Chen Changkang, Phys. Rev. B *60*, 698 (1999).
26. L. B. Ioffe, A. I. Larkin, A. A. Varlamov, and L. Yu, Phys. Rev. B *47*, 8936 (1993).
27. J. C. Campuzano, M. Norman, and M. Randeria, Photoemission in the high-\(T_c\) superconductors in Physics of Superconductors Vol II, edited by K. H. Bennemann and J. B. Ketterson (Springer, Berlin, 2004) p. 167-273.
28. D. van der Marel, Phys. Rev. B *60*, R765 (1999).