Fragile three-dimensionality in the quasi-one-dimensional cuprate PrBa$_2$Cu$_4$O$_8$

A Narduzzo$^1$, A Enayati-Rad$^1$, P J Heard$^2$, S L Kearns$^3$, S Horii$^4$, F F Balakirev$^5$ and N E Hussey$^{1,6}$

1 H. H. Wills Physics Laboratory, University of Bristol, Tyndall Avenue, Bristol BS8 1TL, UK
2 Interface Analysis Centre, Oldbury House, 121 St. Michael’s Hill, Bristol BS2 8BS, UK
3 Department of Earth Sciences, University of Bristol, Wills Memorial Building, Queen’s Road, Bristol BS8 1RJ, UK
4 Department of Applied Chemistry, University of Tokyo, 7-3-1 Hongo, Bunkyo-ku, Tokyo 113-8656, Japan
5 National High Magnetic Field Laboratory, Los Alamos National Laboratory, Los Alamos, NM 87545, USA
E-mail: n.e.hussey@bristol.ac.uk

New Journal of Physics 8 (2006) 172
Received 31 March 2006
Published 1 September 2006
Online at http://www.njp.org/
doi:10.1088/1367-2630/8/9/172

Abstract. In this paper, we report on the experimental realization of dimensional crossover phenomena in the chain compound PrBa$_2$Cu$_4$O$_8$ using temperature, high magnetic fields and disorder as independent tuning parameters. In pure crystals of PrBa$_2$Cu$_4$O$_8$, a highly anisotropic three-dimensional Fermi-liquid state develops at low temperatures. This metallic state is extremely susceptible to disorder, however, and localization rapidly sets in. We show, through quantitative comparison of the relevant energy scales, that this metal/insulator crossover occurs precisely when the scattering rate within the chain exceeds the interchain hopping rate(s), i.e. once carriers become confined to a single conducting element.

$^6$ Author to whom any correspondence should be addressed.
1. Introduction

There is growing evidence that in compounds containing isolated atomic [1] or molecular [2] chains, the conventional Fermi-liquid picture of electron-like quasi-particles fails. How the corresponding one-dimensional state [3, 4] of decoupled spin and charge excitations emerges, however, remains unresolved, prompting the search for compounds whose electronic ground state can be tuned progressively towards one-dimensionality. In metals on the boundary of one-dimensionality, the so-called quasi-one-dimensional (quasi-1D) conductors, the Fermi surface takes the form of pair(s) of parallel corrugated sheets in the plane normal to the conducting chain(s). Provided the two orthogonal interchain hopping energies $2t_i$ (which determine the size of corrugation in each direction) are larger than other relevant perturbations, hopping between chains is coherent and in the absence of charge ordering, a 3D Fermi-liquid ground state is stabilized at low temperature. If this corrugation is ‘smeared out’ and the chains become decoupled, however, theory predicts [3, 4] that even weak interactions will drive the system into the 1D Luttinger-liquid state with its associated phenomenon of spin-charge separation. In order to realize and study the crossover between these two extreme ground states, it is necessary to identify materials where $2t_i$ in both directions is restricted, due to orbital overlap or correlation effects, to energies attainable within the laboratory.

PrBa$_2$Cu$_4$O$_8$ (Pr124) is the non-superconducting analogue of the high-temperature superconductor YBa$_2$Cu$_4$O$_8$ (Y124). Its crystal structure, shown in figure 1, consists of edge-sharing double CuO chain networks (oriented along the crystallographic $b$-axis) sandwiched between sets of CuO$_2$ bilayer plaquettes. Substitution of Pr for Y between the bilayers completely suppresses superconductivity (and mobility) within the CuO$_2$ planes [5] whilst preserving the metallicity of the double chains [6, 7]. This offers a unique opportunity to study the charge dynamics of the cuprate chain in isolation. In this paper, we show that temperature, magnetic fields and disorder can all induce a 3D to 1D crossover in the electronic ground state of Pr124 under laboratory conditions. Whilst dimensional crossover phenomena due to high temperatures [8, 9], high magnetic fields [10]–[12] and even strong correlation effects [13] have been well
documented, to our knowledge, this is the first experimental realization of \textit{disorder}-induced one-dimensionality in a 3D compound.

2. Methods

2.1. Crystal growth and characterization

Single crystals of PrBa$_2$Cu$_4$O$_8$ were grown by a self-flux method in MgO crucibles under high-pressure oxygen gas of 11 atm [7]. The impurity content of three of the crystals used in the disorder study was investigated by means of secondary-ion mass spectrometry (SIMS) as well as electron probe micro-analysis (EPMA). SIMS identified a number of trace impurity elements (including Fe but not Ni or Co) but only three, Mg, Al and Sr, had abundances above the detectability limit of our EPMA measurements (100 ppm). Of these, Mg was by far the most abundant, as expected through contamination with the crucible walls.

2.2. Zero- and pulsed-field resistivity measurements

The resistivities were measured using a standard four-probe ac lock-in technique. For the $\rho_c(T, B)$ measurements, electrical contacts were mounted on the top surfaces of the crystal whilst for the $\rho_a(T, B)$ measurements, they were mounted on the corners. In each case, the other highly resistive direction was shorted out to ensure uniaxial current flow. For in-chain resistivity ($\rho_b$) measurements, the crystal dimensions were recorded using a scanning electron
microscope (SEM). In addition, voltage contacts were placed so as to avoid contamination from the two other highly resistive current directions. An example of the voltage contact configuration for \( \rho_b \) measurements is shown in figure A.1 of the appendix. Due to the smallness of the samples and the finite width of the voltage electrodes, the convention adopted in estimating their distance plays an important role in the determination of the absolute resistivity values. In this study, we have taken the midpoint of the wire (diameter 25 \( \mu m \)) at the sample as the marker for the electrode position. Uncertainties in our estimates of the sample dimensions are 10–15%. In previous studies [12, 14, 15], different markers have been used. As discussed in the appendix, we believe this choice of marker is primarily responsible for the large discrepancies in the absolute magnitudes of the \( \rho_b \) values reported in the literature.

The high-field measurements were performed in the 65 Tesla pulsed magnet at the NHMFL, Los Alamos, USA. A typical 100 ms long magnetic field pulse is produced by discharging 1.6 MJ capacitor bank through a reinforced copper alloy magnet coil. Resistance versus magnetic field is recorded again using a standard high-frequency lock-in technique.

3. Temperature and magnetic field-induced dimensional crossover phenomena in Pr124

It was shown previously [14] that in high-quality Pr124 crystals, electrical resistivity at low \( T \) is metallic in all three orthogonal directions and varies approximately as \( T^2 \), consistent with the development of a 3D Fermi-liquid ground state. The resistive anisotropy at low \( T \) however is extremely large (\( \rho_a : \rho_b : \rho_c(T = 0) \approx 1000 : 1 : 3000 \)) [14], with a similar anisotropy in the ratio of the (squares of the) hopping energies. Moreover, whilst \( \rho_b(T) \) in the purest crystals remains metallic for all \( T < 300 \) K, the interchain resistivities \( \rho_a(T) \) and \( \rho_c(T) \) have maxima around \( T = 150 \) K above which their behaviour becomes thermally activated. These maxima have been interpreted either as a 3D to 1D crossover with increasing temperature or as the emergence of a contribution to \( \rho_{a,c}(T) \) from the insulating CuO\(_2\) planes.

Dimensional crossover phenomena can also be realized in high magnetic fields, due to a field-induced real-space confinement of the charge carriers [16]. In the double-chain cuprate Pr124, the Fermi surface consists of two pairs of corrugated sheets extending normal to the reciprocal space axis \( k_b \). Within a simple tight-binding picture, the \( c \)-axis dispersion (for a single chain) is \( E = -2t^c_\perp \cos(k_c,c) \). For \( B \parallel a \), the dominant Lorentz force \( e[v_F \times B] = ev_FB\hat{c} = \hbar dk_c/dt \) causes carriers to traverse the Fermi sheet along \( k_c \). The sinusoidal corrugation then gives rise to an oscillatory component to the \( c \)-axis velocity \( v^c_\perp = \hbar^{-1}[\partial E^c_\perp / \partial k_c] = [2t^c_\perp / \hbar] \sin(k_c,c) \) and hence to a real-space sinusoidal trajectory with amplitude \( A_c = 2t^{c_\perp} / ev_FB \). Thus \( A_c \) shrinks as \( B \) increases until eventually at \( B^c_{cr} = 2t^{c_\perp} / ev_FB \), \( A_c = c \) and the charge carriers become confined to a single plane of coupled chains. Note that \( B^c_{cr} \) is independent of \( 1/\tau \) and therefore independent of temperature and impurity concentration, as verified experimentally [12, 17]. Due to the quasi-1D nature of the Fermi sheets in Pr124, a similar oscillatory component of amplitude \( A_a = 2t^{a_\perp} / ev_FB \) will also be induced along the \( a \)-axis for \( B \parallel c \). In this case, the crossover field is expected to occur at \( B^c_{cr} = 2t^{a_\perp} / ev_Fa \).

Figure 2(a) shows transverse field sweeps of the \( c \)-axis resistivity \( \rho_c(B) \) (transverse to both the current and the conducting chain) at different fixed temperatures. The inset shows a blow-up of the low-field region (enclosed by a solid rectangle in figure 2(a)), minus the 40 K data for clarity). Below a crossing field \( B^c_{cr} (= 10(1) \) Tesla, shown by a dashed line), \( \rho_c(T) \) is metallic.
Figure 2. (a) Transverse field sweeps of the $c$-axis resistivity $\rho_c(B)$ (transverse to both the current and the conducting chain) at different fixed temperatures. The inset shows a blow-up of the low-field region (enclosed by a solid rectangle in figure 1(a), minus the 40 K data for clarity). The dashed line indicates the crossover field $B_{cr}^c$ (see main text). (b) Transverse field sweeps of the $a$-axis resistivity $\rho_a(B)$. Again the dashed line indicates the crossover field $B_{cr}^a$.

Above $B_{cr}^c$ however, the trend is reversed, implying a metallic/non-metallic crossover in the interchain resistivity as a function of field. (For more details on the precise determination of $B_{cr}^c$, please refer to [12].) As shown in figure 2(b), a similar phenomenon occurs in the reciprocal configuration ($B \parallel c$, $I \parallel a$), though here the corresponding crossover field is extremely high, $B_{cr}^a = 62(2)$ Tesla. This is principally due to the fact the $a \sim c/3$ and therefore a much larger field is required to confine the electrons along the $a$-axis.

From the crossover fields in the transverse $B$ sweeps shown in figures 2(a) and (b), we obtain $2t_{\perp} \sim 45(5)$ K and $2t_{\perp}^a \sim 70(2)$ K. We can now compare these values with estimates of $2t_{\perp}$ from the $T$-dependent resistivity data. Figures 3(a) and (b) show $\rho_c(T)$ and $\rho_a(T)$ data respectively for the same crystals that were used in the magnetic field study. Typically, there are two energy scales that are used as a measure of $2t_{\perp}$ in the interplane(chain) resistivities of low-dimensional metals; the peak in $\rho_{\perp}(T)$ at $T_{max}^{c,a}$ is generally regarded as an upper bound for $2t_{\perp}$ whilst the deviation from the low-$T$ quadratic resistivity gives a lower bound. In Sr$_2$RuO$_4$ for example, the latter criterion has been shown to give very consistent agreement with the value of $2t_{\perp}$ estimated for the most conducting band from quantum oscillation experiments [18]. The insets in figures 3(a) and (b) show $\rho_{c,a}(T)$ versus $T^2$ below $T = 70$ and 110 K respectively. The arrows indicate the temperatures $T_{coh}^{c,a}$ at which $\rho_{c,a}(T)$ first deviates from $T^2$. From these plots we find $T_{coh} = 70(5)$ K, $T_{max} = 130$ K, $T_{coh}^a = 50(5)$ K and $T_{max}^c = 180$ K. Note that the $T_{coh}^{c,a}$ values are in excellent agreement with the values for $2t_{\perp}$ determined from the two sets of magnetic field...
measurements. From this we conclude that $T_{\text{coh}}^{a,c}$ defines the temperature at which $a$- and $c$-axis hopping first begins to lose coherence (i.e. when $k_B T \sim 2t_{\perp}$). The peaks, in contrast, appear to represent the temperature at which all interchain coherence is lost. In the temperature range $T_{\text{coh}}^{a,c} \leq T \leq T_{\text{max}}^{a,c}$ therefore, the chains are very weakly coupled and metallicity is seen to disappear only gradually.

4. Disorder-induced localization in Pr124

We now turn our attention to the intrachain current response. Figure 4(a) shows $\rho_b(T)$ data for four needle-shaped samples $b_1$–$b_4$ taken from the same growth batch. At high $T$, $\rho_b(T)$ is metallic and quasi-linear, the slopes being similar in all four samples suggesting uniform carrier concentration. Only $b_1$ however remains metallic down to the lowest temperature studied ($T = 0.5$ K). For $b_1$, $\rho_b(T) = \rho_{0b} + AT^\alpha$ below 100 K. The coefficient $\alpha = 2.3$ falls within the expected range ($2 < \alpha < 3$) for a quasi-1D Fermi-liquid with dominant electron–electron scattering [19]. The other samples have minima at $T = T_{\text{min}}$ below which $\rho_b(T)$ gradually increases. Figure 4(b) shows a blow-up of the low-$T$ resistivity data for $b_2$, located right at the boundary between the metallic and non-metallic behaviour. The kink in $\rho_b(T)$ at $T_N = 17$ K indicated by an arrow coincides with the antiferromagnetic ordering of the Pr ions [20]. In $b_3$ and $b_4$, this kink manifests itself as a change of slope. Intriguingly, in all crystals that exhibit metallic behaviour down to low $T$ (i.e. those with low residual resistivities, including $b_1$), $\rho_b(T)$ appears unaffected by the Pr ordering.

Several mechanisms for the low $T$ insulating behaviour, including variable range hopping [21], Kondo spin scattering [22], the charge Kondo effect [23], the $\ln(1/T)$ dependence observed in underdoped 2D cuprates [24] and the exponential behaviour expected for a Mott insulator [25],
Figure 4. (a) Zero-field $\rho_b(T)$ data for four single crystals with different levels of disorder. The dashed line overlaying the data of sample $b_4$ is a fit to $\rho_b(T) = AT^{-2/3}$. (b) Blow-up of the low-$T$ $\rho_b(T)$ data for sample $b_2$, on the boundary between metallicity and localization. The dashed line is an extrapolation of the metallic $T^2$ dependence between 45 and 60 K. The arrow indicates the antiferromagnetic ordering temperature $T_N$ of the Pr ions. (c) Blow-up of the low-$T$ $\rho_b(T)$ data for sample $b_3$, highlighting the anomalous $-T$ behaviour below $T_N$ (again marked by an arrow).

were considered but found to be incompatible with $\rho_b(T)$, both above and below $T_N$. Whilst Fe concentrations below 100 ppm can give rise to resistivity upturns [26], we note here that the in-chain magnetoresistance behaviour of all four of these crystals is inconsistent with Kondo scattering [27]. Moreover Mg was the only element whose concentrations were found to scale with the size of the resistivity upturns; $\text{Mg}_{b_2} : \text{Mg}_{b_3} : \text{Mg}_{b_4} = 3500 : 4200 : 4700$ (ppm). We therefore conclude that magnetic impurities were not responsible for the upturns in $\rho_b(T)$.

The large increase in $\rho_b(T)$ of $b_3$ and $b_4$ below $T_{\text{min}}$ argues against weak localization, whilst the very gradual nature of the upturn in $\rho_b(T)$ suggests a lack of charge ordering in Pr124, most probably due to the stabilizing presence of the CuO$_2$ planes. The form of $\rho_b(T)$ for $T_N < T < T_{\text{min}}$ is best represented by a power law, (e.g. dashed line in figure 4(a) for $b_4$ where $\rho_b(T)$ varies as $T^{-2/3}$). Below $T_N$ however, the $T$-dependence of $\rho_b(T)$ changes abruptly in all insulating samples. As illustrated in figure 4(c) for sample $b_3$ for example, $\rho_b(T)$ increases approximately linearly with decreasing $T$ over a decade in temperature between 0.8 and 17 K. The Néel ordering of the Pr ions thus splits the insulating behaviour into two disparate forms, one above $T_N$ which is divergent, and one below which is not, thereby making it difficult to identify the intrinsic transport behaviour in the insulating state.

New Journal of Physics 8 (2006) 172 (http://www.njp.org/)
For a quasi-1D metal, the residual resistivity $\rho_{0b}$ is independent of carrier density. Hence $\rho_{0b}$ can be used to obtain a direct measure of the intrachain mean-free path $\ell$. Given that $b2$ lies on the threshold between metallicity and localization, we can therefore extract an upper bound for the nominal mean-free path $\ell_{cr}$ at the metal/insulator boundary by extrapolating $\rho_{b}(T)$ of $b2$ from high $T$ down to $T = 0$ K. The dashed line in figure 4(b) is an extrapolation of the metallic $T^{2}$ dependence between 45 and 60 K. From this we obtain $\rho_{0b} \leq 8(1) \mu$cm, giving $\ell_{cr} \geq 215(25)$ Å. (Note that there are 2 chains per unit cell in Pr124, and so $\rho_{0b} = \pi a c h / 4 e^{2} \ell$, where $a (=3.88$ Å) and $c (=13.6$ Å) are the $a$- and $c$-axis lattice constants respectively.) Since the $b$-axis lattice constant $b = 3.90$ Å, this is equivalent to more than 50 unit cells. Finally, taking estimates of the Fermi wave vector $k_{F}(=0.2$ Å$^{-1}$) from angle-resolved photoemission spectroscopy (ARPES) [28], we arrive at $k_{F}\ell_{cr} \geq 45(5)$ at the localization threshold. It is important to stress that independent estimates of $\ell_{cr}$ from transverse interchain magnetoresistance measurements that are insensitive to any uncertainties in the crystal dimensions and contact configurations agree very well with the value extracted from $\rho_{0b}$ of sample $b2$, thus supporting the convention for voltage markers adopted here.

Localization at such large values of $k_{F}\ell$ is unprecedented. In 3D Fermi-liquids for example, coherent (Bloch) electron motion is destroyed once $k_{F}\ell < 1$, corresponding to a mean-free path shorter than the de Broglie wavelength. In the normal state of the 2D cuprates, the localization threshold occurs for $k_{F}\ell < 10$ [24], though its origin is as yet unknown. To our knowledge there have been no corresponding experimental studies of the localization threshold in quasi-1D systems. Extensive theoretical studies however have predicted that an insulating phase develops in a strictly 1D system for a vanishingly small amount of random impurities [29]–[32]. The tendency towards localization in sample $b2$ at $k_{F}\ell_{cr} \geq 45(5)$ suggests therefore that a fundamental change in the dimensionality of the electronic system has occurred and indeed, direct comparison of the relevant energy scales confirms this to be the case.

5. Discussion

According to ARPES [28], the Fermi velocity within the chains $v_{F} = 2.5 \times 10^{5}$ ms$^{-1}$. Thus the intrachain scattering rate at the localization threshold $\hbar / \tau_{cr} = \hbar v_{F} / \ell_{cr} \leq 80(10)$ K. As discussed above, $2t_{1}^{a,c}$ (as determined from $B_{cr}^{a,c}$) is also proportional to $v_{F}$. Thus direct comparison of the three energy scales $(2t_{1}^{a}, 2t_{1}^{c}$ and $\hbar / \tau_{cr})$ is in fact independent of the value of $v_{F}$. It is clear therefore that at the localization threshold, $\hbar / \tau$ is comparable to both orthogonal hopping energies. In this circumstance, intrachain scattering becomes sufficient to block coherent wave propagation between chains, making them decoupled electronically; the presence of arbitrary disorder within the (effectively isolated) chains then leads immediately to localization at the lowest temperature, as predicted [29]–[32].

As mentioned in the previous section, the form of $\rho_{b}(T)$ for $T_{N} < T < T_{min}$ is best represented by a power law (e.g., dashed line in figure 4(a)) for $b4$ where $\rho_{b}$ varies as $T^{-2/3}$). Whilst this is consistent with predictions for a Luttinger liquid [25], it should be stressed that the $T$-range is too limited to make any concrete claims. Nevertheless, recent optical [33] and ARPES [34] results do support the emergence of Luttinger-liquid behaviour in Pr124, either in pure Pr124 at high $T$ (i.e. when $k_{B}T > 2t_{1}^{a,c}$) or in Zn-doped Pr124 at low $T$ (presumably when $\hbar / \tau > 2t_{1}^{a,c}$, though there no quantitative estimates of the energy scales were made). From both these studies, the
Luttinger parameter $K_\rho$ is estimated to be $\sim 0.24$, a value compatible [25] with the observation of a quasi-linear $T$-dependence in $\rho_b(T)$ at high $T$, if one assumes a carrier concentration close to, but not quite, $1/4$-filling. Significantly, the resistivity data do not exhibit the exponential behaviour expected for a 1D Mott insulator at low $T$ [25]. The small value of $K_\rho$ suggests that this is due to deviations from commensurability, rather than to the absence of strong interactions.

6. Conclusions

In this paper, we have reported supporting evidence for the first experimental realization of disorder-induced one-dimensionality in a 3D solid at low temperatures. This has only been made possible in Pr124 due to the extremely small values of $2t_\perp (\leq 70 \text{ K})$ in both perpendicular directions. In the quasi-1D organic conductors, such 3D to 1D crossover phenomena are extremely difficult to induce (at low $T$) due to the large coupling in the second direction (in \((\text{TMTSF})_2\text{PF}_6\) for example, $2t_\perp^b \sim 600 \text{ K}$ [13]). This makes Pr124 a rather unique bulk system with which to study physical phenomena on the boundary between Fermi-liquid and Luttinger-liquid ground states that complements existing work on chiral edge states in 2D heterostructures [35]. Whilst it appears that disorder allows the manifestation of one-dimensionality in Pr124 at low $T$, it is, by its very nature, one displaying localized behaviour. Nevertheless, our results highlight one possible route to a delocalized Luttinger-liquid at low $T$, in ultra-pure Pr124 under a tilted (pulsed) magnetic field with components in the $a$–$c$ plane larger than the two crossover fields $B_{\text{cr}}^{a,c}$. Should such field-induced confinement also lead to an insulating state, then the paradigm that all states are localized in real 1D systems will gain yet more empirical confirmation. Finally, let us remark on the observation that the dimensional crossovers in Pr124 appear to occur once the strength of a particular perturbation, e.g. temperature $T$, magnetic field $B$, or scattering rate $1/\tau$, exceeds $2t_\perp$. This contrasts markedly with recent studies on quasi-2D conductors where evidence for interlayer coherence is observed despite the fact that both $k_B T$ and $\hbar/\tau \gg 2t_\perp$ [36]. Clearly the phenomenon of interlayer or interchain decoherence in anisotropic metals is still poorly understood and it is only through further systematic studies that this important issue can be resolved.

Acknowledgments

We acknowledge S Andergassen, C Capan, A Carrington, R Claessen, T Enss, T Giamarchi, S M Hayden, S Kivelson, N Shannon, I Terasaki, V Tripathi and J A Wilson for their helpful comments. We also thank K Nozawa and J Hart for technical assistance. This work was supported by the EPSRC (UK). SH was financially supported by the Leverhulme Fellowship (UK). The work at the NHMFL was supported by the National Science Foundation and DOE Office of Science.

Appendix

We address here the issue of the determination of absolute values of the chain resistivity $\rho_b$. This requires an accurate measurement of crystal dimensions and of distances between contacts. With a high-power optical microscope crystal dimensions can be known to a reasonable degree of

New Journal of Physics 8 (2006) 172 (http://www.njp.org/)
Figure A.1. (a) SEM picture of the pair of voltage contacts for sample $b_2$. The white lines represent different markers used to determine the distance between contacts. (b) A similar picture for sample $b_3$. Using the markers at the mid-point of the Au wires, the distances between voltage pads are 203 and 195 $\mu$m for samples $b_2$ and $b_3$ respectively.

accuracy, 20–25%, and with an SEM, significantly better than that. The chief source of systematic error is the large size of the electrical (voltage) contacts relative to their separation. In order to ensure uniaxial current flow in these highly anisotropic crystals it is necessary to coat conductive paint across the entire sample in both directions perpendicular to the current flow. This invariably leads to large contact pads relative to the sample dimensions. Moreover, it is not obvious which criterion one should use for specifying the distance between the voltage contacts in calculating the resistivity values. It depends to a large extent on the contact resistance of each pad and where within the pad the best electrical contact is made. A convention needs to be adopted by associating a marker with each contact pad. As an example of the importance of the marker scheme, figure A.1 shows two SEM pictures of the voltage contacts for samples $b_2$ and $b_3$.

Choosing the marker for the position of each contact at the Au wire mid-point, at the edge of the Ag paste or at the intermediate point between the two, leads to a factor of two or three difference in the measured distance between contacts, and therefore in the absolute value of $\rho_b$. We believe that the discrepancies between values reported here and those reported previously [12, 14, 15] arise from different conventions adopted to measure the distance between voltage electrodes. In this study, we have used the distance between the mid-point of the Au wires to calculate $\rho_b$.

References

[1] Kim C, Matsuura A Y, Shen Z-X, Motoyama N, Eisaki H, Uchida S, Tohyama T and Maekawa S 1996 Phys. Rev. Lett. 77 4054–7
[2] Claessen R, Sing M, Schwingenschlögl U, Blaha P, Dressel M and Jacobsen C S 2002 Phys. Rev. Lett. 88 096402.1–4
[3] Tomonaga S 1950 Prog. Theor. Phys. 5 544–69
[4] Luttinger J M 1963 J. Math. Phys. 4 1154–62
[5] Horii S et al 1998 Physica C 302 10–22
[6] Terasaki I, Seiji N, Adachi S and Yamauchi H 1996 Phys. Rev. B 54 11993–6
[7] Horii S, Mizutani U, Ikuta H, Yamada Y, Ye J H, Matsushita A, Hussey N E, Takagi H and Hirabayashi I 2000 Phys. Rev. B 61 6327–33
[8] Hussey N E, Kibune M, Nakagawa H, Miura N, Iye Y, Takagi H, Adachi S and Tanabe K 1998 Phys. Rev. Lett. 80 2909–12
[9] Moser J, Gabay M, Auban-Senzier P, Jérome D, Bechgaard K and Fabre J M 1998 Eur. Phys. J. B 1 39–46
[10] Behnia K et al 1995 Phys. Rev. Lett. 74 5272–5

New Journal of Physics 8 (2006) 172 (http://www.njp.org/)
[11] Hussey N E, Mackenzie A P, Cooper J R, Maeno Y, Nishizaki S and Fujita T 1998 Phys. Rev. B 57 5505–11
[12] Hussey N E, McBrien M N, Balicas L, Brooks J S, Horii S and Ikuta H 2002 Phys. Rev. Lett. 89 086601.1–4
[13] Vescoli V, Degiorgi L, Henderson W, Grünner G, Starkey K P and Montgomery L K 1998 Science 281 1181–4
[14] McBrien M N, Hussey N E, Meeson P J, Horii S and Ikuta H 2002 J. Phys. Soc. Japan 71 701–4
[15] Nakada K, Ikuta H, Horii S, Nozaki H, Hirabayashi I and Mizutani U 2001 Physica C 357–360 186–9
[16] Gor’kov L P and Lebed A G 1984 J. Phys. Lett. (France) 45 L433
[17] Waku K, Hussey N E, Sakai F, Nohara M, Takagi H, Nakagawa H, Miura N and Adachi S 2000 J. Phys. Soc. Japan 69 21–4
[18] Mackenzie A P, Julian S R, Diver A J, McMullan G J, Ray M P, Lonzarich G G, Maeno Y, Nishizaki S and Fujita T 1996 Phys. Rev. Lett. 76 3786–9
[19] Oshiyama A, Nakao K and Kamimura H 1978 J. Phys. Soc. Japan 45 1136–46
[20] Li W-H, Wu S Y, Lin Y-C, Lee K C, Lynn J W, Weng S S, Hong I P, Lin J Y and Yang H D 2000 Phys. Rev. B 60 4212–9
[21] Mott N F 1990 Metal–Insulator Transitions (London: Taylor and Francis)
[22] Kondo J 1964 Prog. Theor. Phys. 32 37–49
[23] Matsushita Y, Bluhm H, Geballe T H and Fisher I R 2005 Phys. Rev. Lett. 94 157002.1–4
[24] Ando Y, Boebinger G S, Passner A, Kimura T and Kishio K 1995 Phys. Rev. Lett. 75 4662–5
[25] Giamarchi T 2004 Quantum Physics in One Dimension (New York: Oxford University Press)
[26] Monod P 1967 Phys. Rev. Lett. 19 1113–7
[27] Narduzzo A et al unpublished
[28] Mizokawa T et al 2000 Phys. Rev. Lett. 85 4779–82
[29] Mott N F and Twose W D 1961 Adv. Phys. 10 107–63
[30] Ishii K 1973 Suppl. Prog. Theor. Phys. 53 77–138
[31] Giamarchi T and Schulz H J 1988 Phys. Rev. B 37 325–40
[32] Kane C L and Fisher M 1992 Phys. Rev. Lett. 68 1220–3
[33] Takenaka K, Nakada K, Osuka A, Horii S, Ikuta H, Hirabayashi I, Sugai S and Mizutani U 2000 Phys. Rev. Lett. 85 5428–31
[34] Mizokawa T, Nakada K, Kim C, Shen Z-X, Yoshida T, Fujimori A, Horii S, Yamada Y, Ikuta H and Mizutani U 2002 Phys. Rev. B 65 193101.1–4
[35] Milliken F P, Umbach C P and Webb R A 1996 Solid State Commun. 97 309–13
[36] Singleton J, Goddard P A, Ardavan A, Harrison N, Blundell S J, Schlueter J A and Kini A M 2002 Phys. Rev. Lett. 88 037001.1–4