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Variable range hopping in the spin-glass phase of $La_{2-x}Sr_xCuO_4$

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Abstract

The results of the in-plane transport measurements of strongly underdoped $La_{2-x}Sr_xCuO_4$ films in magnetic fields up to 14 T and at temperatures down to 1.6 K are presented. While at high temperatures the resistivity of the films is metallic-like, the low-T transport is governed by the variable range hopping (VRH) mechanism. Careful analysis shows that the temperature dependence of the pre-exponential factor in Mott's VRH law cannot be neglected and the zero-field conductivity is best described by the formula $\sigma = \sigma_0 T^{-2w} \exp[-(T_0/T)^w]$. In a magnetic field parallel to CuO_2 planes the VRH law is retained, with both parameters, σ_0 and T_0, decreasing by about 10–20% at 14 T. This effect may be related to the decrease of the tunnelling barrier between different antiferromagnetic clusters in the presence of a magnetic field.

1. Introduction

Although the evolution of spin dynamics with hole doping in the antiferromagnetic (AF) background of CuO_2 planes in $La_{2-x}Sr_xCuO_4$ (LSCO) seems to be reasonably well understood [1], new experiments on untwined samples reveal more details of the crystallographic and magnetic structure of this compound. Neutron diffraction experiments show a subtle monoclinic distortion from the orthorhombic symmetry in lightly doped crystals up to $x = 0.03$, caused predominantly by the displacement of the apical oxygen atoms [2]. Outside the AF order, above the Neél temperature, a new field-induced spin ordered state is observed up to $x = 0.01$ [3, 4]. This, in its turn, can reshape our understanding of the normal state transport in LSCO [5, 6], which still offers many intriguing questions. In particular, the role of various interactions which drive the system through the insulator–metal transition in the underdoped regime is still a matter of debate. In strongly correlated systems it may be expected that the electron–electron (e–e) interactions play an important role. The coupling to the magnetic background, the disorder or the possible charge inhomogeneity in the strongly underdoped spin-glass (SG) regime may all give important contributions. The SG regime in LSCO, where frozen AF clusters at low T coexist with metallic-like features of transport at higher T, is especially suitable for studies of this problem.

Angle resolved photoemission spectroscopy (ARPES) shows that at $T = 20$ K quasiparticle (QP) nodal states appear when x reaches 0.03, deep in the SG phase [7]. While the in-plane resistivity at high temperature is a linear function of T, it evolves continuously as T is lowered [8]. Below about 70 K the conductivity of the LSCO becomes proportional to $\ln T$. We have clarified recently that this logarithmic T-dependence is caused by strong e–e interactions, rather than by conventional weak localization in two-dimensional (2D) systems [9, 10]. On further cooling, below the x-dependent temperature which we will call $T_{\ln}$, a gradual change to a variable range hopping (VRH) regime is observed [11–13], where conductivity is described by the relation $\sigma = \sigma_0 \exp[-(T_0/T)^w]$. We focus on this region in the present study. Various values of the exponent w have been found, ranging from 1/4 to 1/2. Usually the temperature dependence of the prefactor σ_0 is neglected, as is done also for other systems [14]. While this gives a reasonable fit to the experimental data in a limited T-range, this may not be in agreement with the theoretical predictions [15]. In particular, the T-dependent prefactor has been predicted for cuprates by a chiral impurity ground state (CIGS) model [16]. It is worth noting that even when the slopes of $\ln\sigma$ versus $T^{-1/4}$ at either end of the T-range depart only by 3% from the mean—and such subtle changes are difficult to observe on standard $\ln\sigma$ versus $T^{-1/4}$ plots—neglecting the prefactor's T-dependence can falsify the estimated value of T_0 by a factor of 2 [17].

Another important issue which still deserves attention is the effect of the magnetic background on the hopping transport.

This problem has been addressed before by the studies of the magnetoresistance (MR) at low T. The MR has been found to be large and negative [9, 10, 18] and correlated with the appearance of magnetic quasi-elastic neutron scattering in the SG phase [19]. The negative MR is believed to be related to a weak-ferromagnetic (WF) transition. This transition, first suggested by the measurements of c-axis resistivity by Thio *et al* [20] in oxygen-doped $La_2CuO_{4+\delta}$, was later found to affect the ab-plane resistivity in LSCO as well [18, 21]. Neutron diffraction directly confirms continuous reorientation of the Cu magnetic moments during the WF transition, from the orthorhombic b-axis to the c-axis when the field is applied along b direction [2]. These findings suggest that the magnetic background has a strong influence on the hopping mechanism. However, there have so far been no studies of how the magnetic field affects the parameters of the VRH law.

In this study we examine the low-T transport and magnetotransport (in the magnetic field parallel to the CuO_2 planes) in LSCO films with x between 0.03 and 0.05. We find that the prefactor σ_0 of the VRH law has a strong T-dependence, but it is inconsistent with the CIGS model. We propose to explain this behaviour by the density of localized states which varies as a power of the energy from the Fermi level, $N(\varepsilon) = N_0\varepsilon^p$, as assumed by Pollak [22] and Hamilton [23], with the small exponent p, of the order of 0.1. This may be interpreted as a result of strong e–e interactions. We also show that both parameters of the VRH law, σ_0 and T_0, decrease in the magnetic field. This result, not anticipated by the existing VRH models, indicates that VRH is sensitive to the state of the magnetic background.

2. Experiment

The c-axis oriented epitaxial films, about 6000 Å thick, were made by pulsed laser deposition on $LaSrAlO_4$ substrates [24]. The characterization and selection of films suitable for the experiments has been described before in detail [9, 25, 26]. Here we describe the conductivity of the films with $x = 0.048$ and 0.045 (one specimen for each x), and for two films, S1 and S2, with the same $x = 0.03$ but with 17% difference in σ at room temperature ($\sigma_{S2} > \sigma_{S1}$). The difference in σ is unlikely to result from the difference in carrier concentration. This is because the laser-ablated films with the same x differ mainly by a small substrate-related strain, while the disorder resulting from the oxygen vacancies plays a minor role [25, 26]. The substrate-related strain leads to the subtle changes of the shape of the Fermi surface, and may, in addition, create small variation in microstructure leading to an increase in grain boundary scattering and the additive term to resistivity. Both effects contribute to a small variation of the transport parameters (including the low-T localization).

The measurements were conducted in the standard four-probe geometry, down to 1.6 K and with the field up to 14 T, applied parallel to the CuO_2 planes. The advantage of using this longitudinal configuration is that it minimizes the influence of the 'classical' Lorentz force. The data below 4.2 K were accumulated by sweeping temperature at constant field, and at higher temperatures mainly by sweeping the field at constant T. Some T-sweeps at higher temperatures were performed as well to verify that both methods give consistent results. We note that performing the T-sweeps at constant B has been instrumental in identifying the location of the low-T anomalies in conductivity related to SG freezing. This is because the temperature of SG freezing in LSCO is B-independent [9], and therefore the field scans at constant T cannot reveal any freezing-related anomalies in the transport.

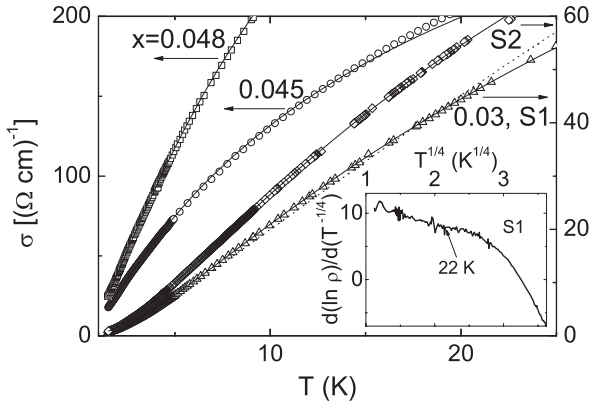

Figure 1. Temperature dependence of LSCO conductivity for $x = 0.048$ and 0.045 (left scale), and two $x = 0.03$ samples, S1 and S2 (right scale). The solid lines are fits of equation (1) to data. The dotted line is the best fit to Mott's VRH formula with a T-independent prefactor for the sample S1, as described in the text. The inset shows the derivative of ln ρ over $T^{-1/4}$ as a function of $T^{1/4}$ for this film.

3. Results and discussion

3.1. Variable range hopping at $B = 0$

Figure 1 shows the conductivity of several LSCO films below T_{ln}, which is equal to about 22, 15 and 10 K for $x = 0.03$, 0.045 and 0.048, respectively. The strong T-dependence of the prefactor in the VRH law is the easiest to observe in films with $x = 0.03$, because in this case the region of the VRH transport extends up to the highest temperature. To show this, we first plot the best fit of Mott's 3D VRH formula with a *T-independent* prefactor, i.e. $\sigma = \sigma_0\exp[-(T_0/T)^{1/4}]$, to S1-sample data. The shape of the fitted line (the dotted line in figure 1) does not mimic the measured σ versus T dependence well and a clear deviation from the data can be seen. In the inset to figure 1 we show the derivative of ln ρ over $T^{-1/4}$ as a function of $T^{1/4}$ for this film. If the prefactor in the VRH law was T-independent then this derivative would be constant. Instead, the derivative increases by about 40% between 22 and 1.6 K and no sign of saturation is observed down to the lowest T.

The variation of the derivative similar to the one shown above is predicted by the theoretical CIGS model [16], in the framework of which some earlier LSCO transport data were re-interpreted. According to this model one expects that the following relation will be fulfilled: $\mathrm{d}(\ln\rho)/\mathrm{d}(T^{-w}) = -AT^w + T_0^w$, with the parameters A and w equal to 6 (4) and 1/4 (1/3) for 3D (2D) hopping, respectively. The fits to our

data below 22 K give instead $A = 3.6 \pm 0.2$ (0.85 ± 0.12) for the 3D (2D) case. Therefore, the variation we observe is too small to be consistent with this model. In another recent model of VRH in cuprates, which assumes a spiral ground state of the magnetic background, the anisotropy of the VRH conduction in LSCO has been calculated [5]. However, it is not clear to us whether a T-dependent prefactor can be explained by this model.

Looking for an alternative description we note that our recent studies indicated the importance of the e–e interactions for the low-T transport behaviour [9, 10]. Therefore, we examine the VRH models which include the effect of the e–e interactions on the density of states. Our data are definitely inconsistent with VRH in the presence of a Coulomb gap which should lead to an exponent $w = 1/2$ [14]. On the other hand, we can fit the conductivity data below $T_{\ln}$ to the formula

$$\sigma = \sigma_0 T^{-2w} \exp[-(T_0/T)^w]. \quad (1)$$

This T-dependence of conductivity in the VRH regime was obtained by Pollak [22] and Hamilton [23] on the assumption that $N(\varepsilon) = N_0\varepsilon^p$ and the exponent w in equation (1) is equal to $(p+1)/(p+4)$. The $p = 0$ limit (i.e. constant density of states) leads equation (1) to Mott's VRH formula with $w = 1/4$, and the $p = 2$ limit leads to VRH in the presence of a Coulomb gap with $w = 1/2$. Equation (1) describes the conductivity of LSCO below $T_{\ln}$ very well. The best fits, shown as solid lines in figure 1, are found for $x = 0.03$ films for non-zero values of p, $p = 0.21 \pm 0.01$ for film S1 (which gives $w = 0.29$) and $p = 0.30 \pm 0.01$ for film S2 ($w = 0.32$). A similar analysis performed for $x = 0.045$ and 0.048 films indicate that the temperature dependence of the prefactor cannot be ignored for these films either. However, the best fits give much smaller values of p (0.05 ± 0.01 and 0.08 ± 0.01, respectively). The rest of the fitting parameters for the most insulating sample, S1, are $T_0 = (6.3 \pm 0.3) \times 10^3$ K and $\sigma_0 = (4.6 \pm 0.2) \times 10^4$ $(\Omega\,\mathrm{cm})^{-1}$. On going from film S1 to $x = 0.048$, i.e. with increase in the the hole concentration, the T_0 decreases by about 70% and σ_0 increases by about 50%, which is consistent with the system becoming more metallic.

Exponent p describes the modification of the density of states $N(\varepsilon)$ in the vicinity of the Fermi level due to the presence of the e–e correlations. Note that p is a result of averaging over the whole momentum space. As such, it is not capable of revealing all details of the electronic structure visible in ARPES experiments [7] and only gives an *effective* $N(\varepsilon)$. The carrier concentration in the two $x = 0.03$ films is most likely very close. The increase in p on going from sample S1 to S2 is accompanied by a 6% increase in the slope of the logarithmic dependence above $T_{\ln}$, $S_{\ln} = \mathrm{d}\sigma/\mathrm{d}(\ln T)$ (caused by e–e interactions [9, 10]). These changes give a consistent picture in which the e–e interactions, modifying the density of states near the Fermi level, also influence the conductivity in the VRH regime at low temperatures.

The exponent p is essentially equal to 0 for the films with $x = 0.045$ and 0.048. This corresponds to Mott's expression for VRH hopping of carriers in a band of localized states in the absence of e–e interactions. This seems to be in apparent contradiction to the fact that ARPES indicates the linear increase of the spectral weight of nodal QP states with increasing x in the SG regime of LSCO [7, 27]. However, in the presence of a sufficiently large carrier concentration one expects that the e–e interactions will be screened, so that the Mott VRH dependence will be recovered as the insulator–metal transition is approached.

3.2. Negative magnetoresistivity

Figure 2(a) shows the MR, defined as $\Delta\rho/\rho_0 = (\rho - \rho_0)/\rho_0$ (where ρ_0 is the resistivity at zero field), for the film S1. A small positive MR (pMR) is observed at high T. As we have demonstrated previously [9, 10], this pMR may be explained by the Zeeman splitting in the particle–hole interaction channel in the disordered electronic system [28]. As the temperature is decreased, a large negative MR (nMR) appears, and grows prominently with decreasing T.

The nMR is frequently observed in spin-disordered systems, for example in canonical spin glasses [29], or in giant magnetoresistance systems [30], where the nMR is believed to be caused by a suppression of the spin-disorder (sd) scattering by the magnetic field. In these systems the sd scattering rate decreases below the SG freezing temperature, T_g, which reduces the magnitude of the nMR. In the Néel state of lightly-doped LSCO a large nMR has been shown to be caused by the reorientation of the WF moments towards the direction of the external magnetic field [20, 21]. The WF moments (WFM), which stem from Dzyaloshinskii–Moryia interactions, are confined to the bc plane [31]. In the AF ordered state at $T < T_N(B)$ and at finite field B applied parallel to the CuO_2 plane, the WFM are rotated from the original orientation at zero field (perpendicular to the CuO_2 planes) by the field-dependent angle $\theta(B)$. This angle becomes equal to $\pi/2$ at $T = T_N(B)$, or at any T when B increases up to its critical value $B_{cr}(T)$. Since at any field below B_{cr}, $\theta(B)$ decreases continuously from $\pi/2$ down to the field-dependent value, $\theta_0(B)$, when temperature decreases from $T = T_N(B)$ down to zero [4], the magnitude of nMR grows continuously with decreasing temperature.

It has been suggested that in the SG region the mechanism of the WF reorientation may be present as well [18]. The onset of the low-T nMR correlates with the growth of neutron scattering in the SG phase [19], so it is clearly caused by a spin-related effect. However, the precise nature of the SG phase in LSCO is still under debate. Some ideas include charge inhomogeneities, or fluctuating stripes, when the charges could self-segregate into the domain boundaries of the AF-correlated clusters [32, 33]. The alternative picture is the one of the noncollinear spiral ground state [5, 34–36], which does not invoke any charge inhomogeneities (which, in fact, have not been observed experimentally so far) and has been capable of explaining some of the peculiarities of magnetism and transport in LSCO. It remains to be seen how any of these scenarios can account for large nMR.

It is instructive to replot the data as magnetic-field-induced changes to conductivity, $-\Delta\sigma \equiv \Delta\rho/\rho_0\rho$, as shown in figure 2(b). This form of plot allows us to emphasize small changes in conductivity induced by the magnetic field, which

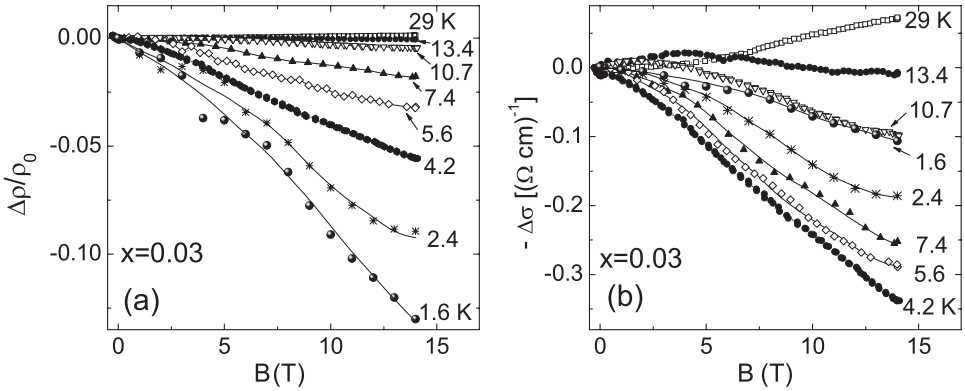

Figure 2. The field dependence of the MR (a) and $-\Delta\sigma$ (b) for fixed temperatures for film S1 with $x = 0.03$. The data for $T < 4.2$ K are calculated from the T-sweeps at constant fields, and therefore are subject to somewhat larger experimental uncertainty. The lines are guides to the eye.

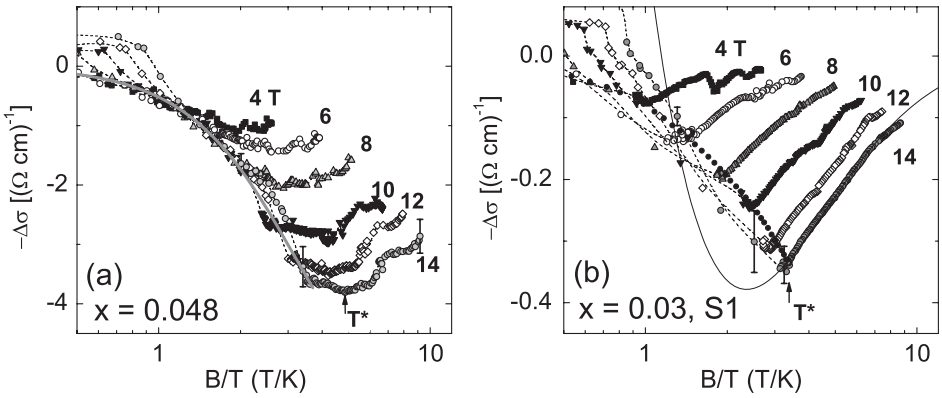

Figure 3. $-\Delta\sigma$ as a function of B/T for films with $x = 0.048$ (a) and $x = 0.03$ (b). Points show the data collected during T-sweeps in constant magnetic fields indicated in the figure. The small black points in (b) show the data accumulated during the B-sweep at $T = 4.2$ K. The thick grey line in (a) shows the common dependence followed by the data, $\tanh^2[c\,(B/T)]$. The black line in (b) is the difference between the fits of equation (1) to $\sigma(T)$ at $B = 14$ and 0 T, as described in the text. The arrows indicate T^* at which $\Delta\sigma$ reaches maximum at $B = 14$ T.

are obscured by the large value of ρ_0 in the original plot of the MR. In contrast to the MR, the $\Delta\sigma$ does not grow monotonically with decreasing T, but has a maximum at low T around 4.2 K, followed by a decrease at still lower T.

In figure 3 we show the dependence of $\Delta\sigma$ on B/T for the film S1, and compare it to a similar plot for the film $x = 0.048$, reproduced from [9][2]. The data were accumulated mostly during the T-sweeps down to 1.6 K, at constant B. In addition, to show the consistency of the results for film S1 (for which the magnitude of $\Delta\sigma$ is small) we plot the data accumulated during B-sweep at $T = 4.2$ K, as small black points in (b). These data points coincide, within experimental error, with the results of the T-sweeps, except for the range of small B/T. In this range $-\Delta\sigma$ measured during the B-sweep is negative whereas the points from the T-sweeps show upward deviations. These upward deviations result from the presence of pMR at high temperatures, which is absent in the B-sweep taken at $T = 4.2$ K.

We note that the magnitude of $\Delta\sigma$ decreases on going from $x = 0.048$ towards $x = 0.03$, by a factor of 2 for $x = 0.045$ film (not shown here, see [10]), and by a factor of 10 for $x = 0.03$ film. This is caused by the reduced number amount of carriers which participate in VRH conduction. Aside from the magnitude, the dependence of $\Delta\sigma$ as a function of B/T is qualitatively the same for all samples, that is, it shows two distinct regimes of behaviour, at high and at low B/T. The temperature T^*, shown in figure 3 for $B = 14$ T, separates these two regimes.

At temperatures $T > T^*$ (i.e. for low B/T), as we have noted previously [9, 10], the $\Delta\sigma$ data, measured in different fields, collapse onto one common curve given by the simple relation $\Delta\sigma \sim \tanh^2[c(B/T)]$. This dependence is shown in figure 3(a) by a thick grey curve. The origin of the common dependence may be qualitatively explained by the spin-related scattering in the SG phase. In AF disordered systems the conductivity is usually found to be proportional to the square of the magnetization [30], and the common curve we find resembles the square of the Brillouin function for $S = 1/2$, but with some effective magnetic moment. The local magnetization in our samples is unknown, but it is possible that it can be roughly approximated in this fashion. The constant c is proportional to an effective magnetic moment and its value

[2] Note that $\Delta\sigma$ was previously calculated in a slightly different way, as $-\Delta\sigma = \Delta\rho/\rho_0^2$.

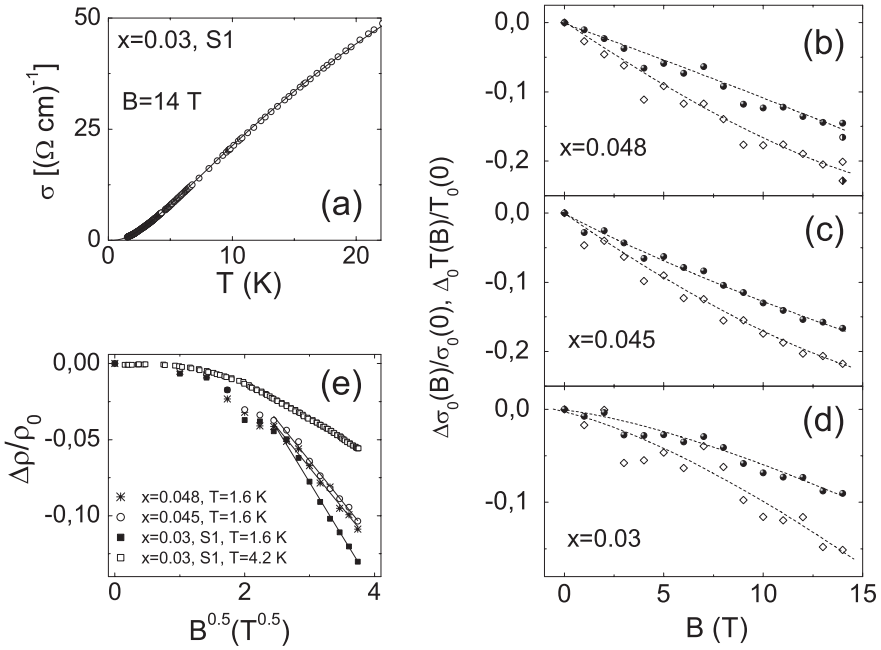

Figure 4. (a) σ versus T measured at $B = 14$ T for film S1. The solid line shows the best fit of equation (1) to data up to 22 K, found for $p = 0.19 \pm 0.01$. (b)–(d) The field-induced change of the parameters of the VRH law, σ_0 (open diamonds) and T_0 (solid circles) for three values of x, found from the fits of equation (1) up to 4.2 K. The lines are guides to the eye. Half-filled symbols for $x = 0.048$ denote results of the fit of equation (1) to $B = 14$ T data up to T^*. (e) Magnetoresistivity at low temperatures as a function of $B^{0.5}$. The solid lines are the best fits to this dependence, starting from 6 T.

is roughly 1.5 times larger for $x = 0.03$ than for $x = 0.048$. This treatment is rather oversimplified and we do not expect that these values of c correspond to a real local moment which contributes to spin scattering. Nevertheless it is interesting to note that the ratio of the c values suggests an increase by a factor of 1.5 of the local moment when x drops from 0.048 to 0.03. Neutron scattering experiments indicate that the frozen moment in the SG phase increases by a factor of about 1.6 in the same x-range [37]. Close correspondence between these ratios supports our conclusions about the origin of a common dependence above T^*.

At $T < T^*$ (high B/T), the data measured at different fields deviate from the common curve. After reaching a maximum at the temperature T^* the magnitude of $\Delta\sigma$ starts to decrease. In case of the film with $x = 0.03$ the maximum is sharp and located at about $T^* \simeq 4.15$ K. Since this location is close to the temperature around which the method of cooling of the cryostat changes (i.e. below 4.2 K the helium is pumped from above the liquid), we have taken extra precautions to rule out the possibility that the maximum is caused by experimental error. The accumulation of the data during B-sweeps, and during T-sweeps allows us to rule out this possibility. In addition, in the films with $x = 0.045$ and 0.048 the maxima occur at lower $T^* \simeq 2.9$ K, and they are definitely less sharp. This gives convincing proof that the feature is real, and it is somehow related to the number of carriers in the film.

There are two possible origins of the maximum in $\Delta\sigma$, both related to the carrier concentration. First, at low T there is a strong localization of carriers, which should be accounted for by the VRH law. In addition, the maximum occurs in the close vicinity of the SG freezing temperature, T_g, which may affect the spin-related scattering. To evaluate the role of these two effects, in the next section we first discuss the effect of the magnetic field on the VRH law.

3.3. VRH in the presence of a magnetic field

Whatever the microscopic mechanism of nMR is, the VRH-like character of $\sigma(T)$ is overall retained in the presence of the magnetic field. This is shown in figure 4(a), where we plot the data taken at $B = 14$ T. The continuous line shows that the data are very well approximated by the VRH law given by equation (1), fitted here up to the temperature $T = 22$ K.

However, we note that the VRH parameters depend slightly on the T-range used for the fit, most likely as a result of the maximum of $\Delta\sigma$ at T^*. Therefore, for a detailed comparison of the influence of magnetic field on the VRH parameters we have restricted the analysis to the data taken below T^* in the case of the $x = 0.03$ film. In the two other films with $x = 0.045$ and 0.048, because the anomaly is broad, it is possible to extend the fitted T-range above T^*, up to 4.2 K. The small T-range used for this analysis causes problems. Namely, the fit with three variable parameters results in a large scatter of their values. Therefore, while we cannot completely rule out the possibility that the exponent p changes with the field, we have decided to keep it constant, and equal to the value obtained from the fit at zero field (carried out in the range covering over a decade in T).

Figures 4(b)–(d) show the magnetic-field-induced relative changes of the two parameters of the VRH law, calculated with respect to their zero-field values, $\sigma_0(0)$ and $T_0(0)$. The changes are defined as $\Delta\sigma_0(B)/\sigma_0(0) = \sigma_0(B)/\sigma_0(0) - 1$,

and $\Delta T_0(B)/T_0(0) = T_0(B)/T_0(0) - 1$. We see that both parameters decrease, with a slightly larger effect for σ_0. The combined effect of both changes is the increase of conductivity in the magnetic field, because the positive change of the exponential term in equation (1) overcomes the negative change of σ_0. The relative changes of the hopping parameters increase with x, and this reflects the increase in the concentration of hopping carriers. In case of $x = 0.048$ we show (by half-filled symbols) the results obtained after cutting the upper limit of the fit down to T^*. This gives an even larger change of σ_0 and T_0 with the field but results in a larger scatter of the fitting parameters for different fields, because of the very limited T-range.

Now we are ready to go back and discuss the origin of the maximum of $\Delta\sigma$ at T^*. Equation (1) is a product of three terms, σ_0, T^{-2w} and the exponential term. The second of these terms, T^{-2w}, is field-independent, and, in addition, it is a monotonic function of T, so it does not contribute to the maximum of $\Delta\sigma$. Using the parameters of the VRH law obtained from the fits, we verify that the maximum may appear on the curve of $\Delta\sigma(T)$ as a result of two opposite tendencies, the decreases of σ_0, and the increase of the exponential term in the field. At low temperatures the increase of the exponential term has a stronger effect on the total conductivity, leading to an increase of $\Delta\sigma$ with increasing T. At high T the decrease of σ_0 becomes more important, and so $\Delta\sigma$, after reaching a maximum, starts to decrease. This effect is illustrated in figure 3(b), where the black solid line depicts the *difference* between the fits of equation (1) to the experimental data measured below 4.2 K in the field $B = 14$ and 0 T.

To show the location and shape of maximum so produced we draw the black line in figure 3(b) well above 4.2 K. We see that the maximum is broad and located at around 6 K. This is in glaring contrast to the actual data in which the maximum in $\Delta\sigma$ is sharp and located at about $T^* \simeq 4.15$ K. In addition, the shape of the black line in the region $T > T^*$ is very different from the experimentally observed dependence $\Delta\sigma \sim \tanh^2[c(B/T)]$. Therefore, the simple extension of the VRH law to higher T cannot fully explain the behaviour of the nMR. Since the maximum around T^* occurs in the close vicinity of the SG freezing temperature, T_g, it is natural to suspect that the spin background plays some role here. Note that, like T_g, T^* increases with decreasing x. Moreover, we expect that the SG freezing should have a more profound effect on the $x = 0.03$ film, in which the AF-correlated clusters are larger. Indeed, the maximum in $\Delta\sigma$ is sharper in the $x = 0.03$ film. It is tempting to try to explain the observed behaviour by a spin-disordered scattering. Naively, one could expect that the strong internal magnetic fields in the SG phase make $\Delta\sigma$ smaller than the value predicted by a VRH law, and indeed this is what seems to be observed. However, in view of the complicated magnetic structure of LSCO this might be a gross oversimplification. Definitely, we can safely conclude that above T^* magnetic and temperature scales are coupled, as is shown by scaling of the $\Delta\sigma$ with the function dependent on B/T. Below T^* these energetic scales are de-coupled. While the temperature T^* is somehow related to the SG freezing, its detailed location is affected by the parameters T_0 and σ_0 describing the VRH hopping in the magnetic field at low temperatures.

3.4. Discussion

The effect of the magnetic field on the VRH parameters has been studied in the past in many materials. In nonmagnetic materials such as doped semiconductors the MR is usually positive, caused by a decrease of the overlap of impurity wavefunctions induced by the magnetic field [14]. In some semiconductors nMR in the VRH regime has been observed in a weak field [38] and explained by the concept of weak localization extended onto the VRH regime [39]. However, this nMR is usually masked by the pMR at higher fields. Interestingly, it has been found in GaAs [40] that the magnetic field which induces this weak-field nMR affects exclusively one hopping parameter, T_0, leaving σ_0 intact.

As shown in panel (e) in figure 4, the negative MR at low temperatures and the highest accessible fields changes linearly with the square root of the magnetic field. This kind of dependence was predicted by one of the models describing nMR in the hopping transport [41–43] and was shown to stem from the change in the localization length proportional to $B^{1/2}$. The nMR in this model arises from the quantum interference (QI) between the direct paths with *forward* scattering (as opposed to the *backscattering* paths in the case of weak localization). The mechanism, originally proposed by Nguen *et al* [41], was next a subject of intense theoretical studies and was used to interpret the nMR in the Mott-type VRH transport, mainly in 2D systems [44–46]. However, the field range where the nMR $\propto B^{1/2}$ dependence is observed in our LSCO films is not consistent with the theory.

According to the numerical calculations [42, 43], the nMR should have $B^{1/2}$ dependence when the field is so strong that the magnetic length, $\lambda = (\hbar/eB)^{1/2}$, is smaller than $(R_h\xi)^{1/2}$, where R_h is the optimum hopping distance. In samples close to the insulator–metal transition the Mott-type VRH behaviour (note: with T-dependent prefactor) is restored ($p \approx 0$; see section 3.1) and R_h can be obtained from the formula $R_h(T) \approx \frac{3}{8}\xi(T_0/T)^{1/4}$, while the localization length ξ is given by the relation $\xi^3 = 49/k_B N_0 T_0$.[3] This, together with the measured density of the states, N_0, being in the range 2–13 states/eV cell [47–49], allow us to estimate roughly that for the sample with $x = 0.048$ ξ is between 13 and 23 Å, and R_h at $T = 1.6$ K is between 34 and 63 Å, correspondingly. This, in its turn, means that the $B^{1/2}$ dependence should be observed according to the theory only in the fields higher than 28 T (or even 156 T for the upper limit of ξ), well outside our measuring window, while in reality it is found to be fulfilled starting from the field ≈6 T. Thus, we cannot conclude that the data are consistent with the theory [42, 43] describing mechanism of QI between the direct paths.

Another mechanism proposed to explain the negative MR in the hopping transport is the field-caused shifting of the levels around the Fermi energy that may enhance the phonon-assisted hopping and thus lead to nMR [50, 51]. This model was developed for the nonmagnetic granular materials where nearest neighbour hopping (NNH), manifesting itself by the $\rho \sim \exp(E/k_B T)$ dependence, is the main

[3] The numerical constant in ξ is based on the percolation parameter equal to 2.7; see [17].

transport mechanism and its predictions agree well with the experimental results on such systems, namely Al/Al_2O_3 [52] and percolating gold films [53]. The magnitude of the predicted nMR is large (up to 80%) and saturates and even starts to decrease with increasing B at the highest fields at low temperatures (see figures 1 and 2 in [51]). According to the theory this nMR happens when the potential barrier of two adjacent grains is small. When the distance (and thus also the potential barrier) between the metallic grains embedded in the dielectric medium increases the MR becomes positive. This is why the negative MR should be observed only near the insulator–metal transition (percolation threshold point in this case) according to this model [50, 51].

The frozen AF clusters in the SG phase of LSCO are far from being similar to the metallic grains in the granular materials. The thermal activation mechanism of transport is absent or at least does not dominate in the temperature region where the nMR is observed. If any analogue of the mechanism of the level-shifting-induced nMR in the nearest neighbour hopping existed in the VRH in LSCO then the magnitude of nMR would decrease with increasing distance from the insulator–metal transition (in the described model [50, 51] the MR even changes its sign to positive). What is observed, however, is increasing magnitude of nMR for the more insulating samples of LSCO. Clearly, the mechanism of nMR must be of another character.

One should remember that both the QI mechanism and the level-shifting-induced nMR described in [50, 51] are of an orbital nature. LSCO is a quasi-2D system and the field during the experiment was applied parallel to the 2D CuO_2 planes with the aim of focusing on the spin-related effects. The exponent $w \simeq 1/4$ found for the $x = 0.045$ and 0.48 films is indicative of 3D hopping and suggests that the hops of carriers in these samples take place not only inside the CuO_2 planes but also between them. Thus, even when the field is applied parallel to these 2D planes, some orbital effects, such as these discussed above, may be observed. Measurements on the untwinned single crystals, with the field rotated in the CuO_2 planes relative to the current, are required to separate the spin and orbital effects better. However, the geometry of the experiment, the quasi-2D character of the samples and the fact that—as we mentioned before—the onset of the low-T nMR correlates with the growth of the magnetic quasi-elastic neutron scattering in the SG phase [19], makes us think that in the measured nMR the spin-related effects are dominant.

The fact that both σ_0 and T_0 in equation (1) change in the field is very intriguing. These VRH parameters depend on the density of the localized states, N_0, and the localization length, ξ, as $\sigma_0 \propto 1/[N_0^{2/(p+4)}\xi^{3(p+6)/(p+4)}]$ and $T_0^{p+1} \propto 1/(N_0\xi^3)$. However, the simple assumption of the field-induced increase of the density of states cannot explain the observed decrease of σ_0 and T_0 with the field. Eliminating field-dependent ξ from the above formulae leads to the proportionality $N_0 \propto \sigma_0/T_0^{(p+1)(p+6)/(p+4)}$ (which, for $p = 0$, reduces to the more familiar Mott dependence $N_0 \propto \sigma_0/T_0^{3/2}$). So estimated, N_0 increases by about 2% for $x = 0.03$ and by 4% for $x = 0.048$ when field increases from zero to 14 T. These reckoned changes are far too small to explain alone the magnitude of the variation of the parameters T_0 and σ_0 observed in experiment (see figure 4). Clearly, another mechanism must be dominant.

There are recent modifications of Mott's original model attempting to explain large nMR in manganites in the paramagnetic phase [54–57]. They introduce the magnetization-dependent probability for an elementary hopping process (and hence modify the probability-related term σ_0) or/and the magnetization-dependent energy barrier for hopping (and thus modify energy barrier related parameter T_0). While the influence of the field on T_0 remains a subject of debate (see [57]), σ_0 is expected to increase with the increase of the local magnetization [56], contrary to the total effect of the field on σ_0 perceived in LSCO.

The incommensurate magnetic order, experimentally observed in LSCO in a broad range of doping [19, 37, 58]—from the Néel phase through the spin-glass phase to the superconducting one—gave rise to the development of the noncollinear spiral ground state model [5, 6, 34–36]. In the framework of this extended t–J model, the staggered component of the Cu moments is represented by a field $\vec{n}$. Besides the term related to the weak ferromagnetism (due to the Dzyaloshinskii–Moriya interaction), two additional terms describing the interaction of the hole with the magnetic field appear in the Hamiltonian. Namely, the magnetic field parallel to the *local* direction of the $\vec{n}$-field increases the projection of the total spin (the sum of spins 'up' and 'down') along $\vec{n}$ and this spin interacts in the Zeeman way with the field. The field perpendicular to $\vec{n}$ tilts the Cu spins allowing the holes to hop more easily between them, what leads to the negative MR.

The influence of the boundaries between AF clusters on the VRH in the spin-glass phase of LSCO has not been discussed yet in the framework of the spiral ground state model. The neutron measurements reveal the presence of two characteristic length scales in the spin-glass state: l_I, related to the incommensurate spin modulation, and the magnetic correlation length, l_M, related to the spin-glass disorder [19, 36, 59, 60]. Since $l_M > l_I$, the role of the boundaries between different clusters must not be disregarded. In a perfectly ordered antiferromagnet the carriers do not change their spin orientation during the hops. The system follows a path of least resistance and the carrier can hop from one ion to the next only if the moments of these ions are not antiparallel [61, 62]. Thus, as pointed out in [56], the hopping barrier should not depend on the local magnetization in the system. In LSCO the dominating long range commensurate AF order directed along the orthorhombic b-axis is destroyed by doped holes at around $x = 0.02$. Even in the Néel phase the mobile hole, when moving in the spiral magnetic background, is forced to change its spin making it parallel to the localized magnetic moment of the background. In the uniformly doped system, the spiral is directed along the crystallographic axis [6]. Intuitively, in the cluster SG phase an additional energetic barrier appears when the carrier crosses the boundary between different clusters. The mechanisms of crossing such boundaries may include both hopping and tunnelling, as it was proposed for manganites [57]. The applied magnetic field should decrease the tunnelling barrier. However, a theoretical evaluation of such a complicated situation is

needed to verify if it may account for the decrease of both VRH parameters as we observe experimentally.

4. Conclusions

In conclusion, we have studied in detail the behaviour of the in-plane transport and magnetotransport in the SG phase of LSCO. In the absence of a magnetic field the temperature dependence of the prefactor in the VRH law cannot be ignored, although it is too small to be consistent with the chiral impurity ground state model. It may be explained by an energy-dependent density of states resulting from the e–e interactions in the VRH regime, previously found to be responsible for $\sigma \propto \ln T$ dependence. The VRH law is retained in the presence of the magnetic field. Both hopping parameters decrease with an increase in the magnetic field, presumably because of the decreased tunnelling barriers at the boundaries of the AF-oriented clusters.

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References

[1] For a condensed review of the recent experimental results see: Julien M-H 2003 *Physica* B **329–333** 693
[2] Reehuis M, Ulrich C, Prokeš K, Gozar A, Blumberg G, Komiya S, Ando Y, Pattison P and Keimer B 2006 *Phys. Rev.* B **73** 144513
[3] Gozar A, Dennis B S, Blumberg G, Komiya S and Ando Y 2004 *Phys. Rev. Lett.* **93** 027001
[4] Neto M B S and Benfatto L 2005 *Phys. Rev.* B **72** 140401(R)
[5] Kotov V N and Sushkov O P 2005 *Phys. Rev.* B **72** 184519
[6] Lüscher A, Misguich G, Milstein A I and Sushkov O P 2006 *Phys. Rev.* B **73** 085122
[7] Yoshida T *et al* 2003 *Phys. Rev. Lett.* **91** 027001
[8] Kastner M A, Birgeneau R J, Shirane G and Endoh Y 1998 *Rev. Mod. Phys.* **70** 897
[9] Cieplak M Z, Malinowski A, Guha S and Berkowski M 2004 *Phys. Rev. Lett.* **92** 187003
[10] Cieplak M Z, Malinowski A, Guha S and Berkowski M 2004 *Physica* C **404** 87
[11] Ellman B, Jaeger H M, Katz D P, Rosenbaum T F, Cooper A S and Espinosa G P 1989 *Phys. Rev.* B **39** 9012
[12] Keimer B *et al* 1992 *Phys. Rev.* B **46** 14034
[13] Chen C Y, Branlund E C, Bae C, Yang K, Kastner M A, Cassanho A and Birgeneau R J 1995 *Phys. Rev.* B **51** 3671
[14] Shklovskii B I and Efros A L 1984 *Electronic Properties of Doped Semiconductors* (New York: Springer)
[15] Castner T G 2000 *Phys. Rev.* B **61** 16596
[16] Lai E and Gooding R J 1998 *Phys. Rev.* B **57** 1498
[17] Mansfield R 1991 *Hoping Transport in Solids* ed M Pollak and B Shklovskii (Amsterdam: Elsevier Science) chapter 10
[18] Ono S, Komiya S, Lavrov A N, Ando Y, Balakirev F F, Betts J B and Boebinger G S 2004 *Phys. Rev.* B **70** 184527
[19] Wakimoto S *et al* 1999 *Phys. Rev.* B **60** R769
[20] Thio T, Thurston T R, Preyer N W, Piocone P J, Kastner M A, Jenssen H P, Gabbe D R, Chen C Y, Birgeneau R J and Aharony A 1988 *Phys. Rev.* B **38** 905
[21] Ando Y, Lavrov A N and Komiya S 2003 *Phys. Rev. Lett.* **90** 247003
[22] Pollak M 1972 *J. Non-Cryst. Solids* **11** 1
[23] Hamilton E M 1972 *Phil. Mag.* **26** 1043
[24] Cieplak M Z, Berkowski M, Guha S, Cheng E, Vagelos A S, RabinowitzD J, Wu B, Trofimov I E and Lindenfeld P 1994 *Appl. Phys. Lett.* **65** 3383
[25] Cieplak M Z, Malinowski A, Karpińska K, Guha S, Krickser A, Kim B, Wu Q, Shang C H, Berkowski M and Lindenfeld P 2002 *Phys. Rev.* B **65** 100504(R)
[26] Malinowski A *et al* 2002 *Phys. Rev.* B **66** 104512
[27] Ino A, Kim C, Nakamura M, Yoshida T, Mizokawa T, Shen Z X, Fujimori A, Kakeshita T, Eisaki H and Uchida S 2000 *Phys. Rev.* B **62** 4137
[28] Lee P A and Ramakrishnan T V 1985 *Rev. Mod. Phys.* **57** 57
[29] Nigam A K and Majumdar A K 1983 *Phys. Rev.* B **27** 495
[30] Jia Y X, Lu L, Khazeni K, Crespi V H, Zettl A and Cohen M L 1995 *Phys. Rev.* B **52** 9147
[31] Lavrov A N, Ando Y, Komiya S and Tsukada I 2001 *Phys. Rev. Lett.* **87** 017007
[32] Zaanen J and Gunnarsson O 1989 *Phys. Rev.* B **40** 7391
[33] Kivelson S A, Bindloss I P, Fradkin E, Oganesyan V, Tranquada J M, Kapitulnik A and Howald C 2003 *Rev. Mod. Phys.* **75** 1201
[34] Hasselmann N, Neto A H C and Smith C M 2004 *Phys. Rev.* B **69** 014424
[35] Juricic V, Benfatto L, Caldeira A O and Smith C M 2004 *Phys. Rev. Lett.* **92** 137202
[36] Sushkov O P and Kotov V N 2005 *Phys. Rev. Lett.* **94** 097005
[37] Matsuda M, Fujita M, Yamada K, Birgeneau R J, Endoh Y and Shirane G 2002 *Phys. Rev.* B **65** 134515
[38] For a review see: Ionow A N and Shlimak I S 1991 *Hoping Transport in Solids* ed M Pollak and B Shklovskii (Amsterdam: Elsevier Science) chapter 12
[39] Altshuler B L, Aronov A G and Khmelnitskii D E 1982 *Pis'ma Zh. Eksp. Teor. Fiz.* **36** 757
[40] Rentzsch R, Friedland K J and Ionov A N 1988 *Phys. Status Solidi* b **146** 199
[41] Nguen V L, Spivak B Z and Shklovski B I 1985 *Zh. Eksp. Teor. Fiz.* **89** 1770
[42] Medina E, Kardar M, Shapir Y and Wang X R 1990 *Phys. Rev. Lett.* **64** 1816
[43] Medina E and Kardar M 1992 *Phys. Rev.* B **46** 9984
[44] Faran O and Ovadyahu Z 1988 *Phys. Rev.* B **38** 5457
[45] Ye Q-y, Shklovskii B I, Zrenner A and Koch F 1990 *Phys. Rev.* B **41** 8477
[46] Arushanov K, Lisunov K G, Malang U, Kloc Ch and Bucher E 1997 *Phys. Rev.* B **56** 1005
[47] Greene L H, Malleta H, Plaskett T S, Bednorz J G and Müller K A 1987 *Solid State Commun.* **63** 379
[48] Uchida S, Takagi H, Hishio K, Kitazawa K, Fueki K and Tanaka S 1987 *Japan. J. Appl. Phys.* **26** L443
[49] Wada N, Muro-oka H, Nakamura Y and Kumagai K 1989 *Physica* C **157** 453
[50] Wang X R and Xie X C 1997 *Europhys. Lett.* **38** 55
[51] Wang X R, Ma S C and Xie X C 1999 *Europhys. Lett.* **45** 368
[52] Pakhomov A B, McLachlant D S, Oblakova I I and Virnik A M 1993 *J. Phys.: Condens. Matter* **5** 5313
[53] Belevtsev B I, Yu B E, Komnik Yu F and Kopeichenko E Yu 1997 *Low Temp. Phys.* **23** 724
[54] Viret M, Ranno L and Coey J M D 1997 *Phys. Rev.* B **55** 8067
[55] Wagner P, Gordon I, Trappeniers L, Vanacken J, Herlach F, Moshchalkov V V and Bruynseraede Y 1998 *Phys. Rev. Lett.* **81** 3980

[56] Yarlagadda S 2000 *Phys. Rev. Lett.* **84** 4017
[57] Wagner P, Gordon I, Trappeniers L, Vanacken J, Herlach F, Moshchalkov V V and Bruynseraede Y 2000 *Phys. Rev. Lett.* **84** 4018
[58] Yamada K *et al* 1998 *Phys. Rev.* B **57** 6165
[59] Matsuda M *et al* 2000 *Phys. Rev.* B **61** 4326
[60] Fujita M, Yamada K, Hiraka H, Gehring P M, Lee S H, Wakimoto S and Shirane G 2005 *Phys. Rev.* B **65** 064505
[61] Zener C 1951 *Phys. Rev.* **82** 403
[62] de Gennes P G 1960 *Phys. Rev.* **118** 141