Dichotomy in the $T$-linear resistivity in hole-doped cuprates

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From analysis of the in-plane resistivity $\rho_{ab}(T)$ of $\text{La}_{2-x}\text{Sr}_x\text{CuO}_4$, we show that normal state transport in overdoped cuprates can be delineated into two regimes in which the electrical resistivity varies approximately linearly with temperature. In the low-temperature limit, the $T$-linear resistivity extends over a very wide doping range, in marked contrast to expectations from conventional quantum critical scenarios. The coefficient of this $T$-linear resistivity scales with the superconducting transition temperature $T_c$, implying that the interaction causing this anomalous scattering is also associated with the superconducting pairing mechanism. At high temperatures, the coefficient of the $T$-linear resistivity is essentially doping independent beyond a critical doping $p_{\text{crit}} = 0.19$ at which the ratio of the two coefficients is maximal. Taking our cue from earlier thermodynamic and photoemission measurements, we conclude that the opening of the normal-state pseudogap at $p_{\text{crit}}$ is driven by the loss of coherence of anti-nodal quasi-particles at low temperatures.

Keywords: superconductors; cuprates; phase diagram; electrical resistivity

1. Introduction

Although the empirical phase diagram of high-$T_c$ cuprates has been drawn countless times, the precise form (i.e. temperature versus doping curve) of the various temperature scales is still a relatively subjective exercise. Only the Néel temperature $T_N$ and $T_c$, defining the onset of long-range antiferromagnetic order and long-range superconducting phase coherence, respectively, are well established experimentally. Other scales are either drawn arbitrarily or follow entirely different trajectories depending on the type of experimental

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probe employed. The ‘coherence’ temperature $T_{\text{coh}}$, often used to describe the emergence of a conventional Fermi-liquid ground state at high doping, is one such example, though by far the most controversial temperature scale is the pseudogap onset temperature $T^*$.

This controversy was articulated neatly in a recent review by Norman et al. [1], in which they summarized the three most common scenarios for $T^*(p)$, with $p$ being the doped-hole concentration. Spectroscopic probes such as angle-resolved photoemission (ARPES), scanning tunnelling microscopy (STM) and Raman scattering find that $T^*$ gradually merges with $T_c$ beyond optimal doping [2]. In this scenario, the pseudogap state is often regarded as a precursor to superconductivity involving strong pairing correlations without long-range phase coherence. Thermodynamic probes such as specific heat and magnetic susceptibility, on the other hand, find a $T^*(p)$ line that cuts the top of the superconducting dome and vanishes at some critical doping level around $p_{\text{crit}} = 0.19$ [3]. This delineation of the two temperature scales is supported by polarized neutron scattering [4] and polar Kerr-effect studies [5], and suggests that the pseudogap state is some kind of ordered state that competes with superconductivity by removing spectral weight, which is never fully recovered upon entering the superconducting phase: $p_{\text{crit}}$ is often then interpreted as a critical end point with (quantum) critical fluctuations associated with this ordered state influencing the physical properties over a large funnel-shaped region in the $(T,p)$ phase diagram in which marked deviations from conventional Landau Fermi-liquid behaviour are manifest.

As most transport coefficients (e.g. resistivity, Hall resistance, thermopower, etc.) vanish at the onset of bulk superconductivity, the outcome of transport probes has been commonly aligned with the third and final scenario of Norman’s review, in which the $T^*(p)$ line simply terminates at the top of the superconducting dome. As such, transport specialists have tended to ‘sit on the fence’ on this issue, though the observation of a broad funnel-shaped region of $T$-linear resistivity has been widely interpreted, in line with similar observations in heavy fermion systems such as YbRh$_2$Si$_2$ [6] in terms of quantum criticality. The location of this putative quantum critical point, however, as well as unambiguous signatures of quantum criticality in the transport behaviour, has remained elusive, largely due to the high upper critical field $H_{c2}$ values in hole-doped cuprates that restrict access to the important limiting low-temperature region below $T_c(p)$.

In a recent report [7], we employed a combination of persistent and pulsed high magnetic fields to expose the normal state of La$_{2-x}$Sr$_x$CuO$_4$ (LSCO) over a wide doping and temperature range, and studied the evolution of $\rho_{ab}(T)$ with carrier density, from the slightly underdoped ($p = 0.15$) to the heavily overdoped ($p = 0.33$) region of the phase diagram. Rather than collapsing to a singular (critical) point, the region of $T$-linear resistivity in LSCO was found to fan out at low temperatures and dominate the low-$T$ response over a very wide doping range. This anomalous or ‘extended’ critical region was wholly unexpected and contrasts markedly with what is observed in YbRh$_2$Si$_2$ [6] and in other candidate quantum critical systems [8]. Our analysis also revealed that the magnitude of the $T$-linear term scaled monotonically with $T_c$ on the strongly overdoped side, but saturated, or was maximal, at a critical doping level $p_{\text{crit}} \approx 0.19$ at which superconductivity

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itself is most robust. The observation of a singular doping concentration in LSCO close to $p = 0.19$ at which a bulk transport property undergoes a fundamental change at low $T$ lends support to the claim that the pseudogap temperature $T^*$ or energy scale $\Delta_g$ vanishes inside the superconducting dome, rather than at its apex.

In this article, we explore this issue further by examining the temperature derivative of the normal-state electrical resistivity of LSCO across the same doping range. By focussing on the temperature range between the zero-field $T_c$ and 500 K, we show that the $T$-linear resistivities at high and low temperatures have different gradients, different doping dependencies and by inference, different origins. In the low-temperature limit, the $T$-linear resistivity appears to result from scattering processes that are associated with the superconducting pairing mechanism. The high-temperature $T$-linear resistivity, by contrast, is attributed to the onset of incoherence, initially of quasi-particles located near the antinodes. At $p = p_{\text{crit}}$, this coherence temperature falls below $T_c$, linking the loss of coherence to the opening of the pseudogap itself.

2. Experiment and results

Single crystals of LSCO with various Sr concentrations were grown using the travelling-solvent floating-zone method and annealed in varying partial pressures of oxygen in order to optimize homogeneity of the oxygen concentration and thus avoid phase separation into hole-rich and hole-poor regions. Crystallographic axes were identified using a Laue camera, and electrical contacts mounted onto individual crystals in such a way as to avoid voltage contamination along the $c$-axis. The low-field ($\mu_0 H < 18$ T) magnetoresistance measurements were carried out using a standard superconducting magnet, while the high-field measurements up to 60 T were performed in a standard $^4$He cryostat at the Laboratoire National des Champs Magnetiques Intenses – Toulouse pulsed-field facility.

Figure 1 shows resistivity data for three representative dopings, $x = 0.17$, 0.21 and 0.26 (denoted LSCO17, LSCO21 and LSCO26, respectively) with the corresponding temperature derivatives plotted in figure 1b,d,f. While the resistivity curves appear smooth and featureless, the derivatives themselves reveal a surprising degree of structure that can be assigned to distinct physical phenomena, each with a characteristic temperature scale. Above $T_c$, $d\rho_{ab}/dT$ rises linearly with increasing temperature (bar the kink at $T_{\text{TO}}$ for the lower doped samples discussed below) until finally it plateaus at a value we label $\alpha_1(\infty)$, the high-temperature limit of the $T$-linear resistivity. Extrapolation of $d\rho_{ab}/dT$ to $T = 0$ results in a second value of the $T$-linear resistivity that we label $\alpha_1(0)$. The reliability of this extrapolation and the recovery of $T$-linear resistivity as $T \to 0$ is highlighted in figure 2 for two different dopings: (a) LSCO23 and (b) LSCO26. In LSCO23, a (pulsed) field of 48 T is sufficient to recover the normal state [7]. The temperature dependence of $\rho(\mu_0 H = 48$ T) is plotted as open squares in figure 2 and compared with estimates (filled squares) of $\rho(0)$, the resistivity extrapolated to zero field, assuming a simple quadratic field dependence of the magnetoresistance. Both exhibit a clear dominant $T$-linear dependence below 60 K. In LSCO26, a field of 18 T is found to extend the range of the $T$-linear resistivity down to at least 5 K, below which residual superconductivity drives

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the resistivity towards zero. Note that the preservation of the $T$-linear resistivity above $H_{c2}$ is counter to recent claims of field-induced quantum criticality via the observation of a $T^2$ $c$-axis resistivity in overdoped $\text{Tl}_2\text{Ba}_2\text{CuO}_6^{+\delta}$ ($\text{Tl}_{2201}$) above $H_{c2}$ [9]. In that particular configuration, $c$-axis magnetoresistance is large and strongly $T$-dependent [10], and, as a consequence, limits one’s ability to extract the intrinsic ($ab$-plane) response.

In our previous paper [7], we showed that for $p > 0.17$, all $\rho_{ab}(T)$ curves for $0 < T \leq 200K$ could be fitted to the expression

$$\rho_{ab}(T) = \rho_0 + \alpha_1(0) T + \alpha_2 T^2,$$  

(2.1)

where $\rho_0$ is the residual resistivity, $\alpha_1(0)$ the coefficient of the low-$T$ $T$-linear term and $\alpha_2$ the coefficient of the quadratic term.

This form of resistivity is self-evident from the temperature derivatives displayed in figure 1, in particular, the linear slope of $d\rho_{ab}/dT$ above the fluctuation regime. A similar expression is also found to describe $\rho_{ab}(T)$ in overdoped $\text{Tl}_{2201}$ at low temperatures [11,12]. Note that the observed form of $d\rho_{ab}/dT(T)$ is inconsistent with the single-component analysis, i.e. $\rho_{ab}(T) = \rho_0 + \alpha_n T^n$, advocated by some groups [13,14] as according to the latter, $d\rho_{ab}/dT(T)$ should exhibit strong curvature and go to zero at $T = 0$. For the highest
Figure 2. (a) Zero-field $\rho_{ab}(T)$ of one LSCO23 crystal (solid line), plotted together with the extrapolated values of $\rho(0)$ (filled squares, see text) and of $\rho(\mu_0 H = 48 T)$ (open squares) for the various temperatures indicated. (b) $\rho_{ab}(T)$ of one LSCO26 single crystal in zero field (0 T) and a magnetic field of 18 T applied parallel to the c-axis. The dashed lines in both (a) and (b) serve to highlight the preservation of the linear $T$-dependence of the resistivity to lower temperatures as superconductivity is destroyed. (Online version in colour.)

For $x$-values ($p \geq 0.24$), $d\rho_{ab}/dT(T)$ does exhibit downward curvature, but only at high temperatures. The latter can be interpreted either as a tendency towards resistivity saturation [7,15,16] or as the development of a non-integer power law in $\rho_{ab}(T)$ at high temperatures, for example, owing to development of ferromagnetic critical fluctuations near the apex of the superconducting dome [17,18].

The form of $\rho_{ab}(T)$ described in equation (2.1) is also consistent with the form of the in-plane transport scattering rate $\Gamma$ extracted from angle-dependent magnetoresistance (ADMR) measurements in overdoped Tl2201 [19,20]. In particular, $\Gamma$ is found to be composed of two components with different $T$ and momentum ($k$) dependencies: one ($\gamma_{iso}$) isotropic and quadratic in $T$, the other ($\gamma_{aniso}$) anisotropic, maximal near the saddle points at $(\pi,0)$ and proportional to temperature. The fact that the two $T$-dependent components in $\Gamma(T,k)$ and $\rho_{ab}(T)$ are additive implies the presence of two distinct, independent quasi-particle

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scattering processes that coexist everywhere on the cuprate Fermi surface. This contrasts with models, e.g. based on hot spots [21] or cold spots [22], in which different regions of \( \mathbf{k} \)-space have different relaxation rates, since in these cases, the lifetimes ought to be additive.

3. Discussion

Figure 3 summarizes the doping dependence of \( \alpha_1(\infty) \) (filled diamonds) and \( \alpha_1(0) \) (filled circles), collated from a large number of LSCO crystals with closely spaced Sr concentrations (including analysis of data for \( 0.15 < x < 0.21 \) reported by [23]). For \( x > 0.30 \), \( \alpha_1(0) \sim 0 \) and the low-\( T \) resistivity is strictly quadratic [24,25]. With decreasing \( x \), \( \alpha_1(0) \) grows rapidly, attaining a maximum value of around \( 1 \mu\Omega \text{cm K}^{-1} \) at \( p_{\text{crit}} = 0.19 \pm 0.01 \). It is interesting to note that this anomalous, distinctly non-Fermi-liquid dependence of the electrical resistivity occurs in a regime of the phase diagram where coherent fermionic quasi-particles, as revealed by the observation of quantum oscillations [26], persist.

The inferred correlation between \( \alpha_1(0) \) and \( T_c \) beyond \( p_{\text{crit}} \) agrees with an earlier ADMR study of \( \Gamma(T, \mathbf{k}) \) in T2201 [27], and has also been seen more recently in the quasi-one-dimensional organic superconductor (tetramethyltetraselenafulvalene)_2X (X = PF_6, ClO_4) [28]. This correlation, coupled with the striking similarity between the angle dependence of \( \gamma_{\text{aniso}} \) and the order parameter symmetry, implies that the interaction responsible for this anisotropic \( T \)-linear scattering rate is also involved in the superconducting pairing. Collectively, these numerous features (the additive components to \( \Gamma(T, \mathbf{k}) \), the form of the anisotropy in \( \gamma_{\text{aniso}} \), the vanishing of the \( T \)-linear component along the nodes, its persistence to low temperatures and across the overdoped regime and its correlation with \( T_c \) all place strong constraints on the development of any theoretical model put forward to explain the pairing mechanism and charge dynamics of hole-doped cuprates.

A \( T \)-linear scattering rate is often indicative of scattering off a bosonic mode. Obvious candidates in the cuprates include phonons, \( d \)-wave pairing fluctuations [22], spin [21] and charge [29] fluctuations. Since all, bar phonons, appear to vanish in heavily overdoped non-superconducting cuprates, however [30,31], it is difficult to single one out at this stage. For a bosonic mode to be the source of \( T \)-linear scattering, the continuation of its linear \( T \)-dependence to very low temperatures is highly constraining, requiring as it does the presence of an extremely low energy scale. A robust \( T \)-linear scattering rate arises, for example, in bipolaron [32] or other two-dimensional boson–fermion mixtures [33,34] due to fermions scattering off density fluctuations of charge 2e bosons, though how this scattering persists beyond \( p_{\text{crit}} \) and the closing of the pseudogap is not yet clear.

An alternative origin for this \( \mathbf{k} \)-space anisotropic scattering is real-space (correlated) inhomogeneity. Indeed, STM measurements on hole-doped cuprates have found evidence for intense anisotropic scattering with a linear energy dependence [35] associated with an inhomogeneous superconducting state. The linear \( T \)-dependence of \( \gamma_{\text{aniso}}(T) \) and its angular dependence are also mirrored in the frequency-dependent single-particle scattering rate \( \Gamma(\omega) \) inferred from ARPES studies on LSCO [36], suggesting that the two derive from the same
origin. This combined linearity in both the temperature and frequency scales is typically referred to as marginal Fermi-liquid (MFL) phenomenology [37], though its preservation over such a wide doping range is inconsistent with models based on conventional quantum criticality. Indeed, problems with the application of such single-parameter scaling hypotheses to the in-plane resistivity of high-$T_c$ cuprates is already well documented [38].

Assuming an effective interaction with the appropriate $d$-wave form factor, the transport scattering rate and electrical resistivity of a two-dimensional metal close to a Pomeranchuk instability was recently shown to follow a $T^{4/3}$ dependence as $T \to 0$ [39,40]. While this gives rise to a scattering rate (and resistivity) with a form approximately similar to that given in equation (2.1) (when combined with impurity scattering and conventional, isotropic electron–electron scattering), there is no implicit correlation between the strength of the $d$-wave scattering and $T_c$ within this model. Finally, in recent functional renormalization group calculations for a two-dimensional Hubbard model, Ossadnik and co-workers [41] have uncovered a strongly angle-dependent $T$-linear scattering term that originates from spin-fluctuation vertex corrections. Significantly, this $T$-linear scattering rate shows strong doping dependence too and vanishes as the superconductivity disappears on the overdoped side, in good agreement with experimental observations [27].

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The high temperature $T$-linear coefficient $\alpha_1(\infty)$ shows a strikingly different doping dependence to $\alpha_1(0)$. Above $p = p_{\text{crit}}$, $\alpha_1(\infty)$ attains a constant value of $\sim 1.0 \pm 0.1 \, \mu\Omega \, \text{cmK}^{-1}$. Such insensitivity to doping leads us to examine whether $\alpha_1(\infty)$ represents some sort of fundamental limit. For a two-dimensional Drude metal,

$$\frac{d\rho_{ab}}{dT} = \frac{2\pi \hbar d}{e^2 v_F k_F} \frac{d(1/\tau)}{dT},$$

(3.1)

where $d$ is the interlayer spacing ($= 0.64 \, \text{nm in LSCO}$), $v_F$ is the Fermi velocity and $k_F$ the Fermi wave vector. Taking typical ($p$-independent) values of $v_F$ ($= 8.0 \times 10^4 \, \text{ms}^{-1}$) and $k_F$ ($= 7 \, \text{nm}^{-1}$) for the anti-nodal states in overdoped LSCO [42], we find that for $\alpha_1(0) = 1 \, \mu\Omega \, \text{cmK}^{-1}$, the momentum-averaged scattering rate $\hbar/\tau \sim \pi k_B T$. Given that the anisotropic scattering rate varies as $\cos^2 2\phi$ within the plane ($\phi$ being the angle between the $k$-vector and the Cu–O–Cu bond direction) [19], this is equivalent to $\hbar/\tau = 2\pi k_B T$ for states near $(\pi, 0)$. Intriguingly, this intense level of scattering corresponds to the so-called ‘Planckian dissipation limit’, beyond which Bloch-wave propagation becomes inhibited, i.e. the quasi-particle states themselves become incoherent [43].

Related to this, ARPES experiments on overdoped cuprates have shown that the onset of $T$-linear resistivity coincides with the loss of in-plane quasi-particle coherence at the anti-nodal points near $(\pi, 0)$, manifest in the disappearance of a peak in the ARPES spectral function [44]. Taken together, these two features imply that the $T$-linear resistivity in overdoped cuprates is a signature of quasi-particle incoherence once the scattering rate near the zone boundary exceeds $2\pi k_B T$. Accordingly, we label the onset of $T$-linear resistivity beyond $p_{\text{crit}}$ in figure 1 as $T_{\text{coh}}$, the onset temperature of anti-nodal coherence (with decreasing temperature).

Below $p = 0.19$, $\alpha_1(\infty)$ starts to rise sharply with decreasing Sr content. Correspondingly, the ratio $\alpha_1(0)/\alpha_1(\infty)$, plotted in figure 3b, maximizes at $p_{\text{crit}}$ at a value close to unity. This maximum is a clear signature of a fundamental change in the quasi-particle response, coincident with the opening of the anisotropic $d$-wave pseudogap as determined by bulk thermodynamic measurements [3]. It is worth noting, however, that this doping level is also the point at which, according to ARPES [42], the Fermi level crosses a van Hove singularity in LSCO. While this may have some impact on the transport properties [45], band-structure calculations have shown that van Hove singularities in cuprates are very system dependent, while the resistivity behaviour is known to be quantitatively universal among the entire cuprate family [16].

Within the present picture, the rise in $\alpha_1(\infty)$ below $p_{\text{crit}}$ is interpreted as a manifestation of the growth of the incoherent region (i.e. as defined here by the condition $\hbar/\tau \geq 2\pi k_B T$) away from $(\pi, 0)$ as the effective interaction continues to intensify with decreasing $x$. As $T$ is lowered, those same incoherent regions become gapped out, leading to a progressive destruction of the large Fermi surface found above $p_{\text{crit}}$ either into Fermi arcs centred along the zone diagonals [46] or into electron and hole pockets [47]. For the anisotropic coefficient $\alpha_1(0)$ however, the reduction in the total number of coherent states below $p_{\text{crit}}$ is more than offset by the removal of the strong scattering sinks near the zone boundary, leading to an overall decrease, or at least a saturation, in $\alpha_1(0)$ with further reduction in carrier density.
Figure 4. Temperature versus doping phase diagram of \( \text{La}_{2-x}\text{Sr}_x\text{CuO}_4 \) as extracted from the temperature derivative of \( \rho_{ab}(T) \). As in figure 1, the labels \( T_{\text{TO}}, T^* \) and \( T_{\text{coh}} \) represent, respectively, the tetragonal to orthorhombic structural transition, the opening of the pseudogap and the onset of quasi-particle incoherence. The dashed lines are all guides to the eye. (Online version in colour.)

The doping dependence of the various temperature scales extracted from \( \frac{d\rho_{ab}}{dT}(T) \) is captured in figure 4. The labels \( T_{\text{TO}}, T^* \) and \( T_{\text{coh}} \) refer to, respectively, the tetragonal to orthorhombic structural transition, the downturn in \( d\rho_{ab}/dT \) for \( p < 0.19 \) and corresponding downturn above \( p_{\text{crit}} \). The effect of the tetragonal to orthorhombic structural transition on \( d\rho_{ab}/dT \) at \( T = T_{\text{TO}} \) is small but well defined (figure 1b). It is manifest similarly in the double-derivative analysis of \( \rho_{ab}(T) \) performed by Ando and co-workers [23] though comparison with actual structural-analysis data [48] informs us that the dashed line depicting \( T^* \) in fig. 2c of their paper is in fact \( T_{\text{TO}} \).

According to the literature, the standard definition of \( T^* \) is the temperature below which \( \rho_{ab}(T) \) starts to deviate from its linear-\( T \) behaviour at high temperature [49]. Different methods for extracting \( T^* \) often lead to markedly different values, though derivative plots are invariably the most sensitive, and therefore return the highest \( T^* \) values [50]. The above definition can be deceiving, however, particularly in LSCO where the deviation in \( d\rho_{ab}/dT \) is upward. As illustrated in figure 1, the form of \( \rho_{ab}(T) \) is remarkably similar on both sides of \( p_{\text{crit}} \), making it difficult to distinguish between pseudogap opening and the coherent/incoherent crossover in that region of the phase diagram. Indeed, the only way these two temperature scales can be distinguished is via their doping dependencies: while \( T^* \) decreases with increasing \( p \), \( T_{\text{coh}} \) shows the opposite trend. Within our experimental uncertainty, it is not yet possible to determine whether \( T^* \) and \( T_{\text{coh}} \) vanish or simply cross around \( p = p_{\text{crit}} \)—a detailed study of more closely spaced doping levels will be required to address this point. Nevertheless, the fact that the ratio \( \alpha_1(0)/\alpha_1(\infty) \) depicted in figure 3b maximizes at a value close to 1 suggests strongly that \( p = 0.19 \) is the point at which both temperature scales vanish. The shaded area in figure 4 serves to reflect this uncertainty.
The location of $T_{\text{coh}}(p)$ for LSCO from our analysis is found to be in excellent agreement with that found for overdoped Bi2212 [44], overdoped Tl2201 and Ca-doped YBa2Cu3O7−δ [14], affirming that it is a generic feature of all overdoped cuprates. Above this line, the anti-nodal quasi-particles are incoherent [44]. In a conventional metal, $\rho(T)$ tends to saturate at high temperatures once the inelastic scattering rate exceeds the effective bandwidth—the so-called Mott–Ioffe–Regel (MIR) limit. In cuprates however, as in other ‘bad’ metals [51], the onset of incoherence is accompanied by a loss of low-frequency spectral weight to frequencies of the order of the on-site coulomb repulsion $U$ [52,53]. This loss of low-frequency spectral weight manifests itself as a dip in the DC conductivity, and thus a rise in the electrical resistivity to values beyond the MIR limit. A connection between this spectral-weight reduction and the preservation of $T$-linear resistivity has yet to be firmly established, though a phenomenological model based on a dominant, anisotropic, $T^2$ scattering rate and proximity to the MIR limit demonstrated a link between the onset of $T$-linearity in cuprates and the loss of coherence at the anti-nodes [15]. Other, more recent proposals attribute the high-temperature $T$-linear resistivity to incoherent transport induced either by scattering off gauge fluctuations [54] or by the motion of hard-core charged bosons [55].

Finally, before concluding, let us consider briefly the quadratic term in $\rho_{ab}(T)$. According to ARPES, the single-particle scattering rate $\Gamma(\omega)$ of overdoped and optimally doped cuprates is also (approximately) quadratic along the nodal directions [56,57]. This combined quadratic temperature and frequency dependence confirm electron–electron (Umklapp) scattering as its origin. As can be seen in the derivative plots of figure 1, however, the $T^2$ term extends in some cases up to $T^*$ or $T_{\text{coh}}$. While this may seem surprising, given that it is a significant fraction of the Fermi energy, it is not uncommon, particularly in correlated transition metal oxides [58].

4. Conclusions

A $T$-linear resistivity, persisting in some cases up to 1000 K [59,60], has been one of the defining characteristics of the normal state of hole-doped cuprates, yet despite sustained theoretical efforts over the past two decades, its origin and its relation to the superconducting mechanism remain a profound, unsolved mystery. In this article, we have analysed the temperature derivative of in-plane resistivity data on a series of closely spaced LSCO single crystals and have uncovered the presence of not one, but two seemingly independent $T$-linear coefficients that persist for all dopings $0.15 \leq x \leq 0.3$. Combining the results summarized in figures 3 and 4, a coherent picture begins to emerge in which the two coefficients reflect different aspects of the normal state, namely, the low-energy effective interaction and quasi-particle decoherence.

For $x > 0.3$, LSCO is a correlated Fermi liquid with a large, $T^2$ resistivity extending over a broad temperature range [24,25]. As the carrier number falls and the system begins to track back towards the Mott insulating state at $p = 0$, an additional anisotropic interaction develops, giving rise to the $T$-linear scattering term. The same interaction also drives up $T_c$, the condensation energy and superfluid density, the latter reaching a maximum at $p = p_{\text{crit}}$ [61]. At this point
however, scattering intensity becomes so strong that those states near \((\pi,0)\)
begin to decohere and given the proportionality with temperature, \textit{decoherence occurs at all finite temperatures}, not just below \(T^*\). With further reduction
in doping, the strength of the interaction continues to rise, causing yet more
states to lose coherence, leading to an overall reduction or saturation in \(\alpha_1(0)\), a
rise in \(\alpha_1(\infty)\) and a corresponding suppression in the condensation energy and
superfluid density.

Within this picture, the strength of the condensate is destroyed below \(p_{\text{crit}}\)
because the interaction that promotes superconductivity ultimately destroys the
very quasi-particles needed to form the condensate. It is tempting to speculate
that the pseudogap also forms in response to this intense scattering, the electronic
ground state lowering its energy through gapping out the incoherent, highly
energetic anti-nodal states and in so doing, preventing scattering of the remnant
quasi-particle states into those same regions. This implies then that, contrary
to much current thinking, the physics behind pseudogap formation is already
prevalent in the strongly overdoped regime, with decoherence, rather than
competing order, being the principal driver. While the two of course may be
intimately linked, the absence of conventional, critical scaling behaviour in both
the thermodynamic and transport properties of cuprates is at odds with the
notion that \(p_{\text{crit}}\) defines some sort of zero-temperature phase transition.

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