

Evidence for spin-fluctuation-mediated superconductivity in electron-doped cuprates

Received: 8 January 2026

Accepted: 18 August 2026

Published online: 19 October 2026

 Check for updates

C. M. Duffy^{1,2,8} , S. J. Tu^{3,8}, Q. H. Chen^{3,4} , J. S. Zhang³, A. Cuoghi¹,
R. D. H. Hinlopen⁵ , T. Sarkar⁶ , R. L. Greene⁶ , K. Jin^{3,4,7}  &
N. E. Hussey^{1,5} 

In conventional, phonon-mediated superconductors, the transition temperature and the linear-in-temperature normal-state scattering rate are linked through the electron–phonon coupling strength. In cuprate high-temperature superconductors, no equivalent relationship between critical temperature and electron–boson coupling has been identified. Here we highlight such a link with electrons scattering off commensurate antiferromagnetic spin fluctuations of variable strength. We use d.c. resistivity and high-field magnetoresistance to extract the scattering rate in electron-doped $\text{La}_{2-x}\text{Ce}_x\text{CuO}_4$ as a function of doping. We reveal a highly anisotropic inelastic component to the scattering rate whose magnitude diminishes markedly as doping increases. The correlation we highlight between the strength of scattering and the critical temperature then identifies low-energy spin fluctuations as the primary pairing glue in electron-doped cuprates. The contrasting magnetotransport behaviour in hole-doped cuprates suggests that the higher critical temperature in those materials cannot be attributed solely to an increase in the coupling strength to the antiferromagnetic fluctuations. The success in modelling $\text{La}_{2-x}\text{Ce}_x\text{CuO}_4$ reinforces the notion that resolving the origin of high-temperature superconductivity in hole-doped cuprates may require more than a simple extension of the Bardeen–Cooper–Schrieffer theory.

The Bardeen–Cooper–Schrieffer theory of superconductivity, where the electron–phonon interaction binds two electrons into a Cooper pair, forms the basis of our present understanding of conventional superconductors. The critical temperature T_c is governed by an electron–phonon coupling constant λ_{ph} through the McMillan formula¹. The same interaction is also known to scatter electrons in the normal (resistive) state. At intermediate temperatures, where the resistivity $\rho(T)$ is T linear, the associated scattering rate $1/\tau = 2\pi\lambda_{\text{tr}}T$. Here λ_{tr} is a measure of the electron–phonon coupling strength as manifested in

charge transport. A near equivalence of λ_{ph} and λ_{tr} has been demonstrated for a wide array of elements and alloys² (Supplementary Fig. 2), affirming the correlation between the pairing strength and the strength of the relevant (transport) scattering process.

Many superconductors of current interest, including heavy fermions, Fe-based superconductors, organic conductors, nickelates and cuprates, are believed to have a pairing mechanism that is non-phononic in origin. A common feature in many of these systems is their proximity to magnetic order, thereby motivating models

A full list of affiliations appears at the end of the paper. ✉ e-mail: caitlin.duffy@lncmi.cnrs.fr; qihongchen@iphy.ac.cn; kuijin@iphy.ac.cn; nigel.hussey@ru.nl

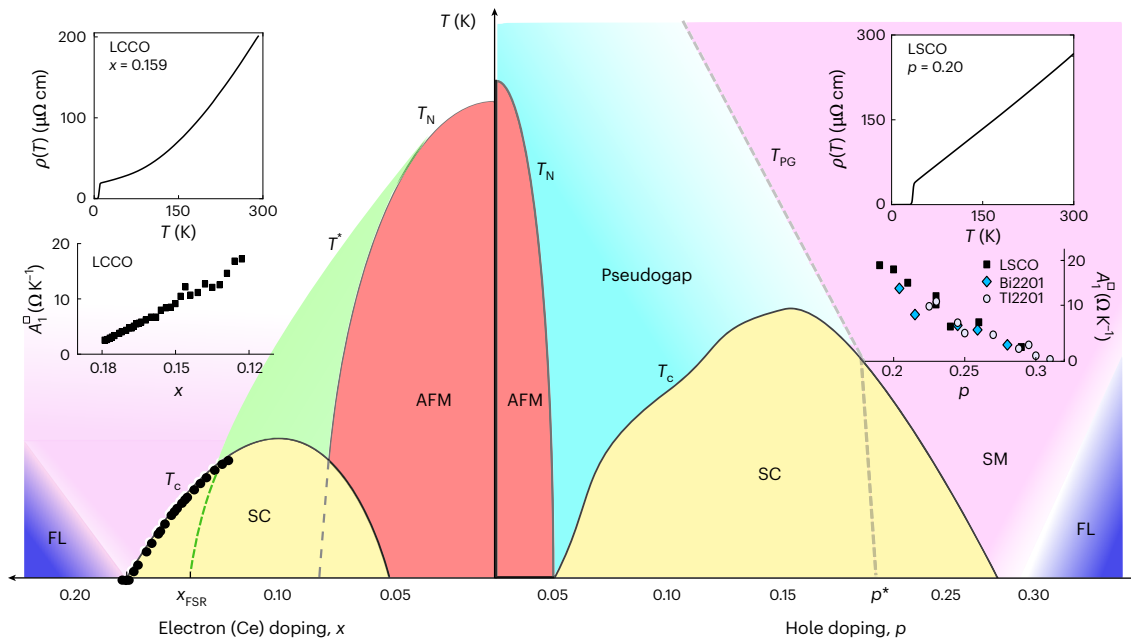


Fig. 1 | Phase diagram of electron- and hole-doped cuprates. On either side of the parent Mott state, the phase diagram of the cuprates exhibits a marked asymmetry. On the hole (p)-doped side, long-range AFM order is suppressed before the onset of superconductivity (SC)³⁹, whereas on the electron (n)-doped side, AFM order coexists with SC although over a narrow doping range. In the latter, short-range AFM correlations persist into the overdoped side (green dome bounded by T^*), causing an FS reconstruction that terminates at x_{FSR} (refs. 40–44). The hole-doped cuprates host SC over a more extended range of doping that peaks at a higher maximum transition temperature. They also play host to the pseudogap phase marked in blue, bounded by T_{PG} . The strange metal (SM) phase lies beyond the pseudogap phase and is characterized by a

predominantly linear-in- T resistivity^{7–10} with a scattering rate corresponding to the Planckian dissipation limit, an example of which is shown in the inset for LSCO at $p = 0.20$. By contrast, the resistivity of the overdoped n -type cuprates ($x > x_{\text{FSR}}$) is approximately T -linear only at the lowest temperatures⁴², becoming quadratic above $T \approx 70$ K (ref. 45). The resistivity of LCCO at $x = 0.159$ is shown in the top left inset. On both sides of the phase diagram, the coefficient of the T -linear resistivity A_1 decreases in tandem with T_c , as shown in the bottom insets for LCCO (from ref. 10), and for three p -type cuprate families (from ref. 13). Both cuprates host a correlated Fermi liquid (FL) phase in the very overdoped, non-superconducting regime.

based on coupling to spin fluctuations^{3,4}. (TMTSF)₂PF₆, for example, undergoes a spin-density-wave transition with an onset temperature that is suppressed with applied pressure, giving way to a dome of superconductivity that encompasses the critical point at which the spin-density-wave phase vanishes. Beyond the ordered phase, a T -linear resistivity is observed at low T whose coefficient A_1 scales with T_c (ref. 5). Weak-coupling theory attributes this correlation to a varying strength of short-range antiferromagnetic (AFM) spin correlations with increasing pressure⁶. So far, however, it has not been possible to experimentally demonstrate a link between the variation in A_1 and a reduction in the inelastic scattering off such low-energy spin fluctuations.

A similar correlation between A_1 and T_c also exists in overdoped cuprates^{7–11}, and by analogy, a common pairing mechanism has been inferred^{18,12}. As with (TMTSF)₂PF₆, however, no link has yet been established between A_1 and a coupling parameter of variable strength. Resistivity is dependent on a number of factors, not simply the strength of scattering, and our inability to isolate the form and strength of $1/\tau$ has led to alternative proposals for the origin of the $A_1 T$ term emerging. It has been argued, for instance, that T -linear scattering in overdoped cuprates is, in fact, universal and tied to the so-called Planckian bound ($\hbar/\tau \approx T$)^{13,14}. In this circumstance, A_1 is determined solely by the (inverse) Fermi temperature T_F (ref. 13) or by the number density of Planckian carriers¹⁴. Even the notion of a T -linear scattering rate in cuprates has been challenged¹⁵. As a result, the identification of the primary pairing mechanism has not been achieved.

In principle, magnetoresistance (MR) measurements can shed light on the origins of the normal-state scattering processes since, according to the Boltzmann transport theory, the fractional MR is independent of the carrier density and is largely determined by the magnitude of $1/\tau$ (extracted from the mean free path ℓ) and its variation

around the Fermi surface (FS). To shed light on the pairing mechanism, it is necessary to demonstrate a correlation between the strength of such scattering and T_c . Although angle-dependent interlayer MR measurements (that is, with $\parallel c$) have had some success in isolating the transport scattering rate in two different families of hole-doped cuprates^{16,17}, corresponding attempts to model the in-plane MR response with the inferred form of $\tau^{-1}(\varphi, T)$ within a Boltzmann approach have so far been met with mixed results^{14,17–20}. No robust connection has yet been made between the in-plane MR response and the nature of the dominant inelastic scattering processes and, thus, between the latter and the strength of superconductivity.

Here we report a detailed and comprehensive study of the in-plane resistivity and MR of the electron-doped cuprate La_{2-x}Ce_xCuO₄ (LCCO) in fine doping steps ($\Delta x \lesssim 0.003$) across the entire overdoped region of the phase diagram, from just above optimal T_c ($x = 0.116$) to the edge of the superconducting dome ($x = 0.188$). The fine doping steps are achieved by a unique combinatorial film growth technique that synthesizes LCCO films with a continuous composition spread under identical conditions on a single SrTiO₃ substrate. More details on the combinatorial technique can be found in ref. 10. The 19 different doping levels investigated in this study are marked in Fig. 1 (solid circles), where the phase diagram for electron- and hole-doped cuprates is presented and the key differences between the two polarities are highlighted.

Our measurements reveal an orbital MR with distinct temperature and doping dependencies that can be captured within a semiclassical Boltzmann treatment by invoking strong anisotropy in the inelastic scattering channel associated with the presence of hotspots. We describe a method to incorporate the effect of current vertex corrections (CVCs) into the Boltzmann equation and apply this to model the Hall coefficient R_H in LCCO, thereby enabling us to identify the location

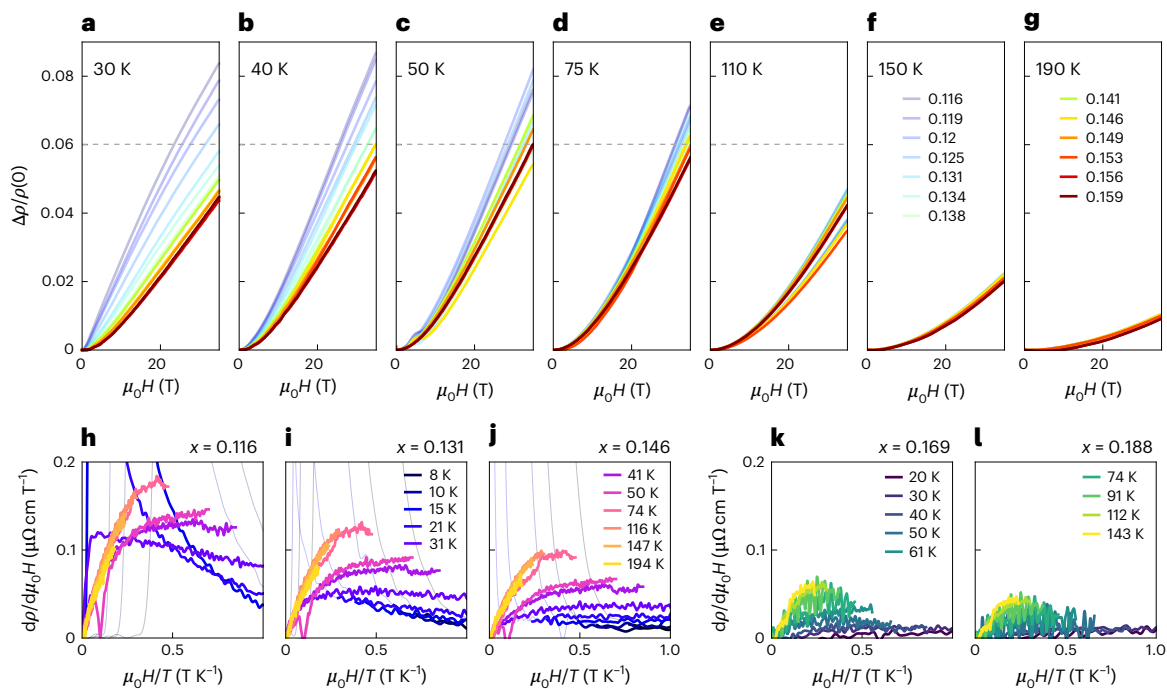


Fig. 2 | T and x dependencies of the in-plane MR in LCCO. **a–g**, MR for $0.116 \leq x \leq 0.159$ on the same film measured across a broad T range ($30 \text{ K} \leq T \leq 190 \text{ K}$). At the highest temperature, the MR is small and quadratic, but as T decreases, the magnitude of the MR grows and becomes more H -linear. For $x > x_{\text{FSR}}$, the MR peaks between 50 and 75 K before falling at lower T . Such behaviour—occurring in the absence of AFM order—implies that the anisotropy in the mean free path ℓ must decrease faster than ℓ itself increases and, thus, must exist predominantly in the inelastic scattering channel (the main text provides more details). The horizontal dashed grey line marks the peak in the magnitude of the MR for $x = 0.159$ at $T = 50 \text{ K}$. **h–l**, Plots of $d\rho/d(\mu_0 H)$ versus H/T for five different x concentrations (three from S1 and two from S2), highlighting

the absence of H/T quadrature scaling across the entire doping range. Below 75 K, H/T scaling fails and the slope of the high-field H -linear MR decreases. This non-adherence to quadrature scaling is in complete contrast to the hole-doped cuprates⁸, suggesting that the origin of their H -linear MR is fundamentally different. At low T , the T dependence of the MR weakens, but the H -linear slope remains. The magnitude of the H -linear MR decreases systematically with increasing x . Similar results on a limited number of $x > x_{\text{FSR}}$ LCCO films were found previously²¹. The thin faint lines in **h–j** represent those sections of the individual MR curves that reside within the mixed state and which are, thus, not included in the subsequent analysis.

of the hotspots on the AFM Brillouin zone (BZ) boundary. The key quantity is the component of this anisotropic scattering rate that scales linearly with temperature and whose magnitude is found to drop by an order of magnitude across the studied doping range. Collectively, this analysis provides evidence that the conduction electrons couple most strongly to commensurate low-energy AFM spin fluctuations in the normal state in LCCO. Given the correlation between the strength of this coupling and T_c , one can then postulate that superconductivity in LCCO and, by extension, in all n -type cuprates, is mediated by AFM spin fluctuations.

Figure 2 summarizes the most salient features of the in-plane MR response in LCCO. Figure 2a–g shows the isotherms of the fractional MR ($\Delta\rho/\rho(0)$, where $\rho(0)$ is the zero-field resistivity) in 13 channels ($0.116 \leq x \leq 0.159$) measured on a single film (S1). The MR is small at high T , grows with decreasing T and x but then peaks at around 50–75 K (for $x \geq 0.12$) before decreasing in magnitude at lower T . Below $x_{\text{FSR}} = 0.14$, a negative component can be seen in the MR at high fields, which has been previously attributed to an isotropic spin effect²¹. The temperature onset of the negative MR component increases with decreasing x as the AFM phase approaches.

In Fig. 2h–l, we plot $d\rho/d\mu_0 H$ —the field derivative of the MR—as a function of H/T at five representative doping levels that span the full studied range. In hole-doped cuprates beyond p^* , all such derivatives (in the normal state) collapse onto a single curve, a form of scaling dubbed Planckian quadrature scaling¹⁸. Although the MR in n -type cuprates exhibits a similar H^2 -to- H crossover with increasing field strength (most evident in Fig. 2j), no such quadrature scaling exists

in LCCO. Moreover, the magnitude of the MR falls monotonically with increasing x (decreasing T_c), despite the fact that the system is becoming more metallic with doping away from the Mott parent state.

As mentioned above, the MR of a metal is set by $1/\tau$ (or more formally ℓ) and its variation around the FS. Typically, this anisotropy is smooth and weakly T dependent, resulting in an MR that obeys Kohler scaling ($\Delta\rho \propto f(H/\rho) \propto f(\omega_c \tau) \propto f(\ell)$, where ω_c is the cyclotron frequency). With decreasing T , τ increases as ρ decreases. Hence, one expects the MR to increase in magnitude as T decreases, as seen at elevated temperatures²². In LCCO, however, the MR below 60 K (for $x > x_{\text{FSR}}$) is found to decrease at a much faster rate than ρ itself decreases. Qualitatively, this observation implies that the size of the MR in LCCO is primarily governed by the degree of anisotropy in the inelastic scattering rate. The drop in $\Delta\rho$ across the film (at fixed T) then further reveals that this anisotropic inelastic scattering weakens as T_c decreases.

To put this on a more quantitative footing, we model the MR using Boltzmann transport theory. Since the Fermi velocity v_F has only modest ($\leq 10\%$) angular variation²³, we can presume that all T -dependent anisotropy in ℓ is contained within $\tau^{-1}(\varphi)$. To model $\rho(H, T)$ successfully, the form of $\tau^{-1}(\varphi)$ should reproduce the following key features of the data:

1. As shown in Extended Data Fig. 1 and Supplementary Fig. 1, the zero-field resistivity below 60 K can be described accurately by the form $\rho_0 + A_1 T + A_2 T^2$. With increasing x , A_1 falls whereas A_2 remains constant, suggesting that only the T -linear term is x dependent¹⁰.

- Since the magnitude of the MR falls with decreasing T , the anisotropy in ℓ must lie predominantly in the inelastic (T -dependent) channels.
- To account for the robust H -linear MR seen at high field strengths (in some cases, up to 65 T (ref. 24)) within the same single-fluid Boltzmann picture, we also need to introduce intense scattering centres (hotspots) that act to impede cyclotron motion^{19,25} and thereby preserve the H -linear form of the MR.

We proceed by applying the Shockley–Chambers tube integral formalism (SCTIF) to the LCCO FS depicted in Fig. 3 (details of this formalism are provided in the Supplementary Information, along with a discussion of the viability of including inelastic scattering processes within a Boltzmann approach). Treating the MR as an orbital effect is justified by the fact that the longitudinal MR is much smaller than the transverse MR (Extended Data Fig. 2).

Our ansatz for the scattering rate $\tau^{-1}(\varphi, T, x)$ is then constructed as follows. To satisfy point (1) above, we introduce three terms in $\tau^{-1}(T, x)$: an impurity term τ_{imp}^{-1} (assumed to be isotropic in accord with the quasi-isotropic ν_F in LCCO), a quadratic term $\alpha_2 T^2$ (the magnitudes of both τ_{imp}^{-1} and α_2 are fixed throughout) and a doping-dependent T -linear term $g(x)\alpha_1 T$. To account for (3)—the H -linear MR—we also introduce a hotspot term $\tau_{\text{HS}}^{-1}(\varphi)$ that comprises a series of sharp peaks (strong scattering centres) whose locations depend on whether the dominant scattering is from spin or charge fluctuations (Fig. 3).

Although these four components (τ_{imp}^{-1} , $g(x)\alpha_1 T$, $\alpha_2 T^2$ and $\tau_{\text{HS}}^{-1}(\varphi)$) can account for $\rho(0, T)$ and the crossover to H -linear MR at high fields, they fail to capture the anomalous growth in the MR with increasing T shown in Fig. 2 (we have assumed here that $\tau_{\text{HS}}^{-1}(\varphi)$ has no intrinsic T dependence). To address point (2), therefore, we must also incorporate φ dependence into one or both of the T -dependent terms of $\rho(0, T)$, that is, $g(x)\alpha_1 T$ and $\alpha_2 T^2$ such that as $\tau^{-1}(T)$ increases, so does its anisotropy. In the following, for simplicity, we assign it to both, that is,

$$\tau^{-1}(\varphi, T, x) = \overbrace{\tau_{\text{imp}}^{-1} + \tau_{\text{HS}}^{-1}(\varphi)}^{T\text{-independent}} + \overbrace{(g(x)\alpha_1 T + \alpha_2 T^2)\beta(\varphi)}^{T\text{-dependent}} \quad (1)$$

Equation (1) contains the minimal ingredients for $\tau^{-1}(\varphi, T, x)$ to reproduce both $\rho(0, T)$ and the T and H dependencies of the MR of LCCO across the doping series ($T > T^*$). The variation in $k_F(\varphi)$ and $\nu_F(\varphi)$ (Fig. 4a(i)) are consistent with the tight-binding parameterization²⁶ presented in Extended Data Table 1. Given that $\nu_F(\varphi)$ is essentially isotropic, the most obvious and intuitive way to incorporate this anisotropy is to associate it with a broadening of the scattering rate around the hotspots themselves, which is, in turn, a measure of the inverse correlation length of the spin or charge correlations. The parameterizations of $\tau_{\text{HS}}^{-1}(\varphi)$ and $\beta(\varphi)$ are somewhat arbitrary and, as extensive analysis revealed, have only a modest influence on the quality of the fitting (other than to set the departure from H -linear MR). Nevertheless, the most successful fits occurred when the maxima of $\beta(\varphi)$ and $\tau_{\text{HS}}^{-1}(\varphi)$ coincided. For illustration, Fig. 4a(ii)–(iv) shows the MR data for $x = 0.159$ and in Fig. 4b,c, the corresponding simulations for two distinct forms of $\tau^{-1}(\varphi, T, x)$, based, respectively, on dominant spin (here $\beta(\varphi) = \sin^{12}[2(\varphi \pm (\pi/4 - \text{HS}))]$ with $\text{HS} = 0.48$) and charge ($\beta(\varphi) = \cos^2[2\varphi]$) scattering.

It is clear that both scenarios reproduce essentially all of the features of $\rho(H, T)$, including the absolute magnitude. Hence, from analysis of MR alone, it is not possible to deduce whether spin or charge fluctuations are the dominant source of anisotropic scattering in LCCO (for completeness, alternative forms of anisotropy that fail to account for key aspects of the MR are presented in Extended Data Fig. 3, whereas a comparison of simulations with different parameterizations for $\beta(\varphi)$ is presented in Supplementary Fig. 3). In order to identify the dominant scatterer, we need to consider the unusual T and x dependencies of the Hall coefficient R_H in LCCO^{10,27,28}.

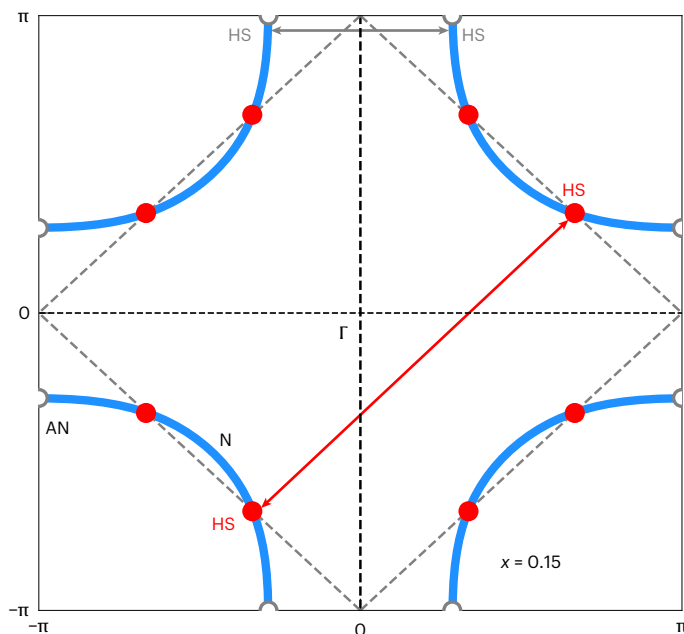


Fig. 3 | Fermi surface of LCCO ($x = 0.15$). The thick blue line represents the in-plane FS. The long diagonal (red) arrow represents large $Q(\pi, \pi)$ hotspot (HS) scattering across opposite sides of the AFM BZ boundary (marked by the dashed grey line). This is consistent with earlier neutron scattering and photoemission studies, showing that spin correlations in electron-doped cuprates remain commensurate within the SC doping range^{46,47} and the corresponding hotspots at the intersection between the FS and the AFM zone boundary⁴⁸. The short horizontal (grey) arrow represents $(\delta, 0)$ HS scattering across flat sections of the FS, presumably induced through charge correlations^{49–53}. The locations of the antinodes (AN) and nodes (N) are also marked. Note that there are eight hotspots in the former case, and only four in the latter.

In a layered system, the Hall conductivity σ_{xy} is determined by the Stokes area swept out by the $\vec{\ell}$ vector as φ varies around the FS²⁹. The FS of LCCO (Fig. 3) has a singular curvature, meaning that $\vec{\nu}_F$ has the same circulation at all φ . In this circumstance, even the inclusion of strong anisotropy in $\tau^{-1}(\varphi)$ cannot cause σ_{xy} (or R_H) to change sign. In LCCO ($x > x_{\text{FSR}}$), however, $R_H(T)$ drops monotonically from its isotropic (Drude) value with increasing T and eventually becomes negative (Extended Data Fig. 4). In the absence of any additional FS reconstruction, for which there is no reported evidence, the only effective way to explain the sign change in $R_H(T)$ is a breakdown of the relaxation time approximation and the incorporation of current vertex corrections into the calculation for σ_{xy} (ref. 28).

In nearly AFM metals, scattering between points on the FS displaced by the dominant scattering vector $\vec{Q}$ leads to mixing of (the velocity of) those states. When the scattering is sufficiently strong as to form hotspots, the total current $\vec{j}_k$ at the hotspots is approximately parallel to $\vec{\nu}_k + \vec{\nu}_{k+Q}$ and is, thus, no longer perpendicular to the FS. Applying linear response theory, Kontani demonstrated how $\vec{j}_k$ varies in proximity to the AFM hotspot³⁰. Following Kontani's model, we incorporate the effect of CVCs into our Boltzmann analysis and recalculate σ_{xx} , σ_{xy} and R_H (Methods and Extended Data Fig. 5). As shown in Extended Data Fig. 6, using the same form for $\tau^{-1}(\varphi)$ that was used to model the MR, we are able to capture the key features of $R_H(T, x)$ beyond $x > x_{\text{FSR}}$, most notably, the reduction in R_H with increasing T . Moreover, CVCs arising from AFM spin fluctuations have been shown to generate a $T + T^2$ resistivity³⁰, consistent with what is found in LCCO (Extended Data Fig. 1).

A similar CVC analysis was carried out for the second scenario in which hotspot scattering is due to charge fluctuations. In this case, however, we find that mixing states across the flat sections of FS does

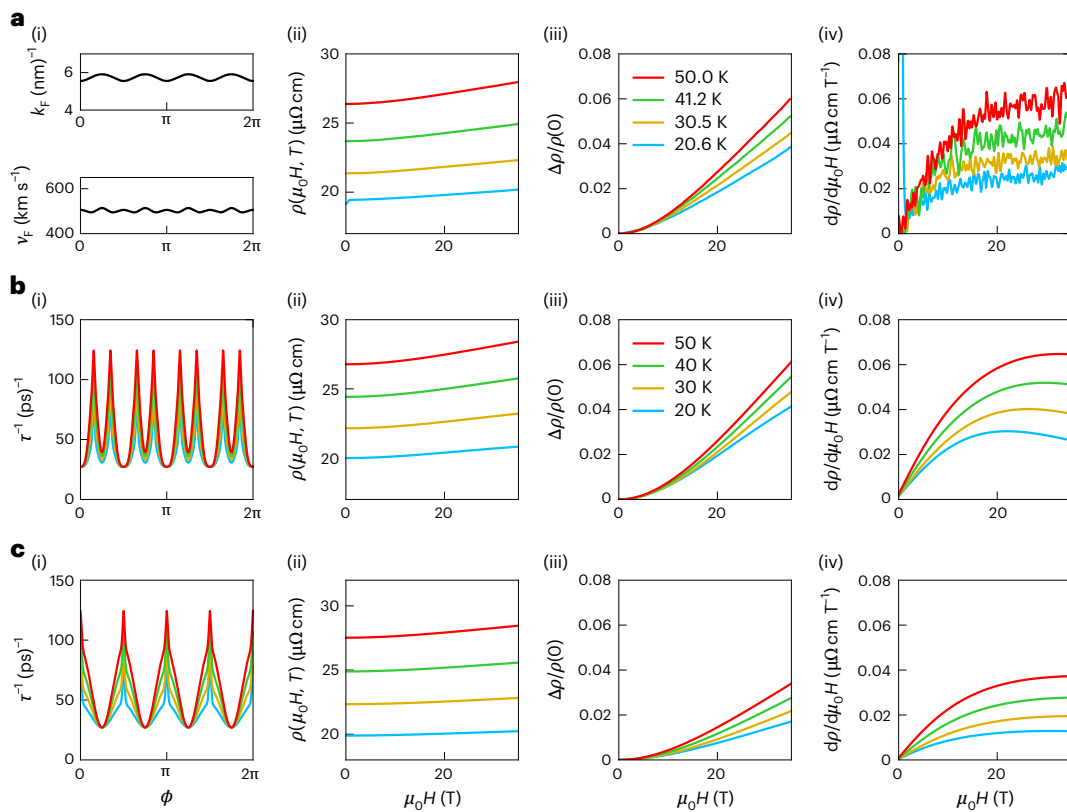


Fig. 4 | Exemplar simulations of MR for $x = 0.159$ ($20 \text{ K} \leq T \leq 50 \text{ K}$).

a, Variation in the FS (top) and Fermi velocity v_F (i) obtained from the tight-binding parameterization given in Extended Data Table 1. Experimental MR curves (ii), $\Delta\rho/\rho(0)$ (iii) and $d\rho/d\mu_0H$ (iv) for $x = 0.159$. Note that the magnitude of the MR increases with increasing T , indicative of developing anisotropy in the inelastic scattering channel, and that the MR undergoes a clear H^2 -to- H crossover, most evident in (iv). **b**, (i)–(iv) Simulations resulting from the form of scattering

rate expressed in equation (1) and depicted in (i), where the inelastic component peaks at the hotspots located at the AFM BZ boundary. The parameters for this simulation are listed in Extended Data Table 2. **c**, (i)–(iv) Corresponding simulations for the form of scattering depicted in (i), with peaks at the flat sections of FS (antinodes) connected by short- Q charge correlations. Both scattering rates capture all the main features of the experimental data, as shown in **b**(ii)–(iv) and **c**(ii)–(iv).

not give rise to a change in circulation of $\vec{J}_k$ ($\vec{v}_k$); hence, as the anisotropic inelastic scattering grows with increasing T , $R_H(T)$ can only become more positive, in marked contrast to what is observed experimentally (Extended Data Fig. 6). Further support for dominant (π, π) scattering is provided through the inspection of the high-field slope of the MR over the full (T, x) range (Extended Data Fig. 7). This is the first key result of this study.

To examine the doping dependence $\tau^{-1}(x)$, we proceed to fit the absolute values of the zero-field resistivity (at $T = 40 \text{ K}$, that is, above T_{FSR}) and the MR (at $T = 40 \text{ K}$ and $\mu_0H = 20 \text{ T}$) for all (~ 13) voltage pairs measured on film S1, using equation (1) – with AFM hotspot scattering – and only one doping-dependent variable parameter – $g(x)$. The results are summarized in Fig. 5 and the corresponding fitting parameters and $g(x)$ values are listed in Extended Data Table 3 for both samples. Figure 5a,b show the threefold drop in $\rho(0, 40 \text{ K})$ and the twofold drop in $\Delta\rho/\rho(0)$ (at 20 T) across S1, respectively. For the MR to obey Kohler scaling as a function of x , one expects $\Delta\rho/\rho(0)$ to increase by a similar amount to the decrease in $\rho(0)$. The counter-trend in $\Delta\rho/\rho(0)$ (Fig. 5b), however, inevitably requires an even greater decrease in $g(x)$ (Fig. 5c). Coupled with the results from S2, the order-of-magnitude drop in $g(x)$ across the series mirrors the variation in A_1 (the coefficient of the T -linear resistivity)¹⁰ depicted in Fig. 1. Over the same doping range, T_c is reduced from near its optimal value (25 K) to zero. This clear correlation between T_c and the strength of the anisotropic T -linear scattering rate is the second key result of this study, revealing the effective coupling parameter for superconductivity in LCCO. These considerations also reveal why the MR is so constraining: (i) in isolating the contribution

of ℓ to $\rho(T)$, (ii) in revealing the presence of marked anisotropy in $\tau^{-1}(\varphi)$ (as well as the presence of hotspots through the observation of a robust H -linear MR) and (iii) in revealing the marked drop in the strength of this anisotropic scattering across the doping series (and with it, the drop in the effective coupling strength λ). Taken together with the identification of (π, π) scattering as the dominant normal-state scatterer, we conclude that SC pairing in electron-doped cuprates is mediated by the exchange of low-energy spin fluctuations.

Before closing, there are five additional points to make. First, it has been argued previously that the magnitude of T -linear resistivity in cuprates is consistent with the notion of a Planckian bound on scattering¹³. The analysis presented here, however, reveals that the magnitude of the scattering rate in LCCO varies markedly across the doping series. Hence, the low- TT -linear resistivity in electron-doped cuprates cannot be tied to the Planckian limit. Second, our study also reveals that the T -linear scattering rate in cuprates is not always isotropic¹⁷. Third, the persistence of the T -linear term in $\rho(T)$ down to low T supports the arguments of Rosch that in the presence of impurity scattering, any shorting effects arising from the colder regions of the FS away from the hotspots are diminished³¹. Additionally, given the similarities between the phase diagrams of n -type cuprates and the Bechgaard salts and the scaling between A_1 and T_c observed in both families^{5,10}, it is perhaps timely to investigate the MR response of the latter to elucidate $\tau^{-1}(\varphi, T)$ and to determine whether an equivalent coupling parameter could also be identified there.

This brings us finally to the dichotomy between the electron-doped cuprates and their hole-doped counterparts. The asymmetry in their

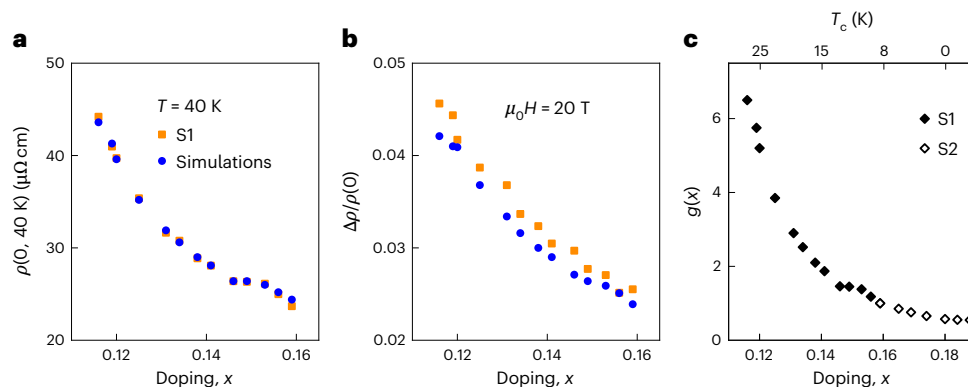


Fig. 5 | Coupling parameter for high- T_c superconductivity in electron-doped cuprates. **a**, Doping-dependent simulations (blue circles) of $\rho(0, 40\text{ K})$ —the zero-field resistivity at $T = 40\text{ K}$ —where signatures of AFM order are absent or weak across the doping range of film S1. The orange squares represent the experimental data and the closed blue circles, the simulations. **b**, Corresponding simulations (blue circles) of the fractional MR $\Delta\rho/\rho(0)$ at $\mu_0H = 20\text{ T}$ and $T = 40\text{ K}$ as a function of x . **c**, Scaling factor $g(x)$ of the anisotropic T -linear scattering

rate as a function of x . Note that $g(x)$ was the only doping-dependent variable parameter used to obtain the results in **a** and **b**. Note, too, that for S2, we take $g(x)$ from the coefficient of T -linear resistivity, and normalize it to the value obtained in S1 at $x = 0.159$. Across the doping series, $g(x)$ drops by one order of magnitude as $T_c \rightarrow 0$. Thus, $g(x)$ directly relates the anisotropic T -linear scattering rate to T_c , acting as an effective coupling parameter for both transport and superconductivity.

respective phase diagrams (Fig. 1) has often led to the assertion that doping the parent Mott insulator with electrons or holes leads to markedly different low-energy physics. Certainly, the p-type cuprates are more strongly correlated³², lying within the strong-coupling regime where U/t is about twice as large as in the electron-type cuprates (U is the on-site Coulomb repulsion and t is the nearest-neighbour hopping parameter). The ratio t/J , where $J = 4t^2/U$ is the AFM exchange interaction, is also larger in the hole-doped cuprates, indicating a lower propensity to remain AFM³³.

Given these differences, it might seem surprising that the low- T resistivities in both wings exhibit such similar behaviour, namely, $\rho_{ab}(T) = \rho_0 + A_1T + O(T^2)$ with a coefficient A_1 that scales with T_c . Our present study has shown that such scaling in LCCO reflects directly the coupling to low-energy AFM spin fluctuations and that weak-coupling calculations exploring this link^{6,34,35} may well reach consensus with experiment if applied to the electron-doped side. It is naturally tempting to speculate that the same correspondence is at play in p-type cuprates and that the higher T_c values in the latter are nothing more than a consequence of a larger λ_{sf} or comparable λ_{sf} coupled with larger U . If that were the case, one would expect the MR response of overdoped p-type cuprates to exhibit qualitatively similar behaviour, that is, one governed by an anisotropic inelastic scattering rate whose magnitude would drop with decreasing T . Yet, across the entire strange metal regime in hole-doped cuprates, the MR response is fundamentally different, in the sense that the MR exhibits a form of quadrature H/T scaling that no variant of kinetic quasiparticle theory has yet been able to account for^{14,18}. Instead, there is growing experimental evidence that the strange metallic state on the hole-doped side has a dual character, playing host to both coherent (Boltzmann-like) and incoherent (non-Boltzmann-like) carriers^{14,18,36}. Turning this around, one might argue that the inability to apply Boltzmann theory to the hole-doped strange metal is a reflection of the fact that the correct parameterization has not yet been found. Its wholesale applicability to LCCO, however, provides a strong counter-argument that such a parameterization should already have been found, were it to exist. This inability to extrapolate our findings in LCCO to the hole-doped side gets to the heart of the problem (the difference between the electron- and hole-doped cuprates) and implies that the higher T_c values in the latter cannot be simply due to an enhanced coupling to spin fluctuations. This also suggests that an understanding of both the strange metallicity and (high-temperature) superconductivity in hole-doped cuprates is likely to require a fundamentally new approach, one that lies beyond the quasiparticle picture^{37,38} or an extension of the Bardeen–Cooper–Schrieffer theory to incorporate

spin-fluctuation-mediated pairing. For the latter, at least, it would appear that electron-doped cuprates represent its epitome.

Online content

Any methods, additional references, Nature Portfolio reporting summaries, source data, extended data, supplementary information, acknowledgements, peer review information; details of author contributions and competing interests; and statements of data and code availability are available at <https://doi.org/10.1038/s41567-026-03454-6>.

References

- McMillan, W. L. Transition temperature of strong-coupled superconductors. *Phys. Rev.* **167**, 331–344 (1968).
- Allen, P. B. in *Handbook of Superconductivity* (ed. Poole, C. P.) (Academic Press, 2000).
- Scalapino, D. J., Loh, E. & Hirsch, J. E. d -wave pairing near a spin-density-wave instability. *Phys. Rev. B* **34**, 8190–8192 (1986).
- Chubukov, A. V., Pines, D. & Schmalian, J. *Superconductivity: Conventional and Unconventional Superconductors* (eds. Bennemann, K. H. & Ketterson, J. B.) (Springer, 2008).
- Doiron-Leyraud, N. et al. Correlation between linear resistivity and T_c in the Bechgaard salts and the pnictide superconductor $\text{Ba}(\text{Fe}_{1-x}\text{Co}_x)_2\text{As}_2$. *Phys. Rev. B* **80**, 214531 (2009).
- Bourbonnais, C. & Sedeki, A. Superconductivity and antiferromagnetism as interfering orders in organic conductors. *C. R. Phys.* **12**, 532–541 (2011).
- Cooper, R. A. et al. Anomalous criticality in the electrical resistivity of $\text{La}_{2-x}\text{Sr}_x\text{CuO}_4$. *Science* **323**, 603–607 (2009).
- Jin, K. et al. Link between spin fluctuations and electron pairing in copper oxide superconductors. *Nature* **476**, 73–75 (2011).
- Putzke, C. et al. Reduced Hall carrier density in the overdoped strange metal regime of cuprate superconductors. *Nat. Phys.* **17**, 826–831 (2021).
- Yuan, J. et al. Scaling of the strange-metal scattering in unconventional superconductors. *Nature* **602**, 431–436 (2022).
- Harada, K. et al. Revised phase diagram of the high- T_c cuprate superconductor $\text{Pb-doped Bi}_2\text{Sr}_2\text{CaCu}_2\text{O}_{8+\delta}$ revealed by anisotropic transport measurements. *Phys. Rev. B* **105**, 085131 (2022).
- Taillefer, L. Scattering and pairing in cuprate superconductors. *Annu. Rev. Condens. Matter Phys.* **1**, 51–70 (2010).
- Legros, A. et al. Universal T -linear resistivity and Planckian dissipation in overdoped cuprates. *Nat. Phys.* **15**, 142–147 (2019).

14. Ayres, J. et al. Universal correlation between H -linear magnetoresistance and T -linear resistivity in high-temperature superconductors. *Nat. Commun.* **15**, 8406 (2024).
15. Pelc, D. et al. Resistivity phase diagram of cuprates revisited. *Phys. Rev. B* **102**, 075114 (2020).
16. Abdel-Jawad, M. et al. Anisotropic scattering and anomalous normal-state transport in a high-temperature superconductor. *Nat. Phys.* **2**, 821–825 (2006).
17. Grissonnanche, G. et al. Linear-in temperature resistivity from an isotropic Planckian scattering rate. *Nature* **595**, 667–672 (2021).
18. Ayres, J. et al. Incoherent transport across the strange-metal regime of overdoped cuprates. *Nature* **595**, 661–666 (2021).
19. Hinlopen, R. D. H., Hinlopen, F. A., Ayres, J. & Hussey, N. E. B^2 to B -linear magnetoresistance due to impeded orbital motion. *Phys. Rev. Res.* **4**, 033195 (2022).
20. Ataei, A. et al. Electrons with Planckian scattering obey standard orbital motion in a magnetic field. *Nat. Phys.* **18**, 1420–1424 (2022).
21. Poniatowski, N. R., Sarkar, T. & Greene, R. L. Anomalous normal-state magnetotransport in an electron-doped cuprate. *Phys. Rev. B* **103**, 125102 (2021).
22. Dagan, Y. et al. Origin of the anomalous low temperature upturn in the resistivity of the electron-doped cuprate superconductors. *Phys. Rev. Lett.* **94**, 057005 (2005).
23. Horio, M. et al. Suppression of the antiferromagnetic pseudogap in the electron-doped high-temperature superconductor by protect annealing. *Nat. Commun.* **7**, 10567 (2016).
24. Sarkar, T. et al. Correlation between scale-invariant normal-state resistivity and superconductivity in an electron-doped cuprate. *Sci. Adv.* **5**, eaav6753 (2019).
25. Koshelev, A. E. Magnetotransport of multiple-band nearly antiferromagnetic metals due to hot-spot scattering. *Phys. Rev. B* **94**, 125154 (2016).
26. Tang, C. Y. et al. Suppression of antiferromagnetic order in the electron-doped cuprate $T'-\text{La}_{2-x}\text{Ce}_x\text{CuO}_{4+\delta}$. *Phys. Rev. B* **104**, 155125 (2021).
27. Sarkar, T. et al. Fermi surface reconstruction and anomalous low-temperature resistivity in electron-doped $\text{La}_{2-x}\text{Ce}_x\text{CuO}_4$. *Phys. Rev. B* **96**, 155449 (2017).
28. Jenkins, G. S. et al. Origin of the anomalous Hall effect in the overdoped n -type superconductor $\text{Pr}_{2-x}\text{Ce}_x\text{CuO}_4$: current-vertex corrections due to antiferromagnetic fluctuations. *Phys. Rev. B* **81**, 024508 (2010).
29. Ong, N. P. Geometric interpretation of the weak-field Hall conductivity in two-dimensional metals with arbitrary Fermi surface. *Phys. Rev. B* **43**, 193–201 (1991).
30. Kontani, H. Anomalous transport phenomena in Fermi liquids with strong magnetic fluctuations. *Rep. Prog. Phys.* **71**, 026501 (2008).
31. Rosch, A. Interplay of disorder and spin fluctuations in the resistivity near a quantum critical point. *Phys. Rev. Lett.* **82**, 4280 (1999).
32. Weber, C., Haule, K. & Kotliar, G. Strength of correlations in electron- and hole-doped cuprates. *Nat. Phys.* **6**, 574–578 (2010).
33. Macridin, A., Jarrell, M., Maier, T. & Sawatzky, G. A. Physics of cuprates with the two-band Hubbard model: the validity of the one-band Hubbard model. *Phys. Rev. B* **71**, 134527 (2005).
34. Moriya, T. & Ueda, K. Spin fluctuations and high temperature superconductivity. *Adv. Phys.* **49**, 555–606 (2010).
35. Wu, W., Wang, X. & Tremblay, A.-M. S. Non-Fermi liquid phase and linear-in-temperature scattering rate in overdoped two-dimensional Hubbard model. *Proc. Natl Acad. Sci. USA* **119**, 2115819119 (2022).
36. Chang, J. et al. Anisotropic breakdown of Fermi-liquid quasiparticle excitations in overdoped $\text{La}_{2-x}\text{Sr}_x\text{CuO}_4$. *Nat. Commun.* **4**, 2559 (2013).
37. Phillips, P. W., Hussey, N. E. & Abbamonte, P. Stranger than metals. *Science* **377**, eabh4273 (2022).
38. Patel, A. A., Guo, H., Esterlis, I. & Sachdev, S. Universal theory of strange metals from spatially random interactions. *Science* **381**, 790–793 (2023).
39. Ando, Y., Lavrov, A. N., Komiya, S., Segawa, K. & Sun, X. F. Mobility of the doped holes and the antiferromagnetic correlations in underdoped high- T_c cuprates. *Phys. Rev. Lett.* **87**, 017001 (2001).
40. Yu, W., Higgins, J. S., Bach, P. & Greene, R. L. Transport evidence of a magnetic quantum phase transition in electron-doped high-temperature superconductors. *Phys. Rev. B* **76**, 020503 (2007).
41. Charpentier, S. et al. Antiferromagnetic fluctuations and the Hall effect of electron-doped cuprates: possibility of a quantum phase transition at underdoping. *Phys. Rev. B* **81**, 104509 (2010).
42. Fournier, P. et al. Insulator-metal crossover near optimal doping in $\text{Pr}_{2-x}\text{Ce}_x\text{CuO}_4$: anomalous normal-state low temperature resistivity. *Phys. Rev. Lett.* **81**, 4720–4723 (1998).
43. Thurston, T. R. et al. Antiferromagnetic spin correlations in $(\text{Nd}, \text{Pr})_{2-x}\text{Ce}_x\text{CuO}_4$. *Phys. Rev. Lett.* **65**, 263–266 (1990).
44. He, J. et al. Fermi surface reconstruction in electron-doped cuprates without antiferromagnetic long-range order. *Proc. Natl Acad. Sci. USA* **116**, 3449–3453 (2019).
45. Sarkar, T., Greene, R. L. & Das Sarma, S. Anomalous normal-state resistivity in superconducting $\text{La}_{2-x}\text{Ce}_x\text{CuO}_4$: Fermi liquid or strange metal? *Phys. Rev. B* **98**, 224503 (2018).
46. Fujita, M. et al. Low-energy spin fluctuations in the ground states of electron-doped $\text{Pr}_{1-x}\text{LaCe}_x\text{CuO}_{4-\delta}$ cuprate superconductors. *Phys. Rev. Lett.* **101**, 107003 (2008).
47. Yamada, K. et al. Neutron scattering study on electron-hole doping symmetry of high- T_c superconductivity. *J. Phys. Chem. Solids* **60**, 1025–1030 (1999).
48. Armitage, N. P. et al. Anomalous electronic structure and pseudogap effects in $\text{Nd}_{1.85}\text{Ce}_{0.15}\text{CuO}_4$. *Phys. Rev. Lett.* **87**, 147003 (2001).
49. Li, Y. et al. Hole pocket-driven superconductivity and its universal features in the electron-doped cuprates. *Sci. Adv.* **5**, eaap7349 (2019).
50. da Silva Neto, E. H. et al. Doping-dependent charge order correlations in electron-doped cuprates. *Sci. Adv.* **2**, e1600782 (2016).
51. da Silva Neto, E. H. et al. Charge ordering in the electron-doped superconductor $\text{Nd}_{2-x}\text{Ce}_x\text{CuO}_4$. *Science* **347**, 282–285 (2015).
52. Ishii, K. et al. High-energy spin and charge excitations in electron-doped copper oxide superconductors. *Nat. Commun.* **5**, 3714 (2014).
53. Hepting, M. et al. Three-dimensional collective charge excitations in electron-doped copper oxide superconductors. *Nature* **563**, 374–378 (2018).

Publisher's note Springer Nature remains neutral with regard to jurisdictional claims in published maps and institutional affiliations.

Open Access This article is licensed under a Creative Commons Attribution 4.0 International License, which permits use, sharing, adaptation, distribution and reproduction in any medium or format, as long as you give appropriate credit to the original author(s) and the source, provide a link to the Creative Commons licence, and indicate

if changes were made. The images or other third party material in this article are included in the article's Creative Commons licence, unless indicated otherwise in a credit line to the material. If material is not included in the article's Creative Commons licence and your intended use is not permitted by statutory regulation or exceeds the permitted

use, you will need to obtain permission directly from the copyright holder. To view a copy of this licence, visit <http://creativecommons.org/licenses/by/4.0/>.

© The Author(s) 2026

¹High Field Magnet Laboratory (HFML-FELIX) and Institute for Molecules and Materials, Radboud University, Nijmegen, Netherlands. ²INSA-T, Université Grenoble Alpes, CNRS, LNCMI-EMFL, Université de Toulouse, Toulouse, France. ³Beijing National Laboratory for Condensed Matter Physics, Institute of Physics, Chinese Academy of Sciences, Beijing, China. ⁴Songshan Lake Materials Laboratory, Dongguan, China. ⁵H. H. Wills Physics Laboratory, University of Bristol, Bristol, UK. ⁶Maryland Quantum Materials Centre and Department of Physics, University of Maryland, College Park, MD, USA. ⁷Key Laboratory of Vacuum Physics, School of Physical Sciences, University of Chinese Academy of Sciences, Beijing, China. ⁸These authors contributed equally: C. M. Duffy, S. J. Tu. ✉e-mail: caitlin.duffy@lncmi.cnrs.fr; qihongchen@iphy.ac.cn; kuijin@iphy.ac.cn; nigel.hussey@ru.nl

Methods

Zero-field resistivity

Combinatorial LCCO thin films were grown epitaxially on SrTiO₃ substrates from targets with $x = 0.10$ and 0.19 , facilitating measurements over a large region of the LCCO phase diagram in minute steps in x . A comprehensive description of the deposition and patterning processes can be found elsewhere (ref. 10). Two films, S1 ($0.116 \leq x \leq 0.159$) and S2 ($0.159 \leq x \leq 0.188$), were patterned into microbridges and their in-plane resistivity measured using a standard four-point a.c. lock-in method. Seven channels could be measured simultaneously on a single film and in total, 13 were measured on S1 and 7 on S2 with one common doping level at $x = 0.159$. As shown in Extended Data Fig. 1a, the two $\rho(T)$ curves at $x = 0.159$ coincide below $T = 70$ K provided that the data for S2 are scaled (reduced) by a factor of 1.5. This same scaling factor was then applied to all channels in S2. Note that the data for film S1 (Extended Data Fig. 1a) and Supplementary Fig. 1 are only discrete points (the zero-field values for each field sweep) rather than a continuous curve. To obtain a full, uninterrupted $\rho(0, T)$ curve between 2 K and 290 K inside the Bitter magnet would have required at least 8 h of continuous measurement, time that was not available at the high-field facility when measuring S1. Fortunately, for S2, the cryostat was available for an extended period following the measurement, making it possible to perform a warming run in zero-field over 3 days.

Extended Data Fig. 1b,c shows the zero-field in-plane resistivity $\rho(0, T)$ and the corresponding derivatives $d\rho/dT$ for a selection of channels on S2, respectively. As clear from Extended Data Fig. 1e, $\rho(0, T) \propto A_1 T + A_2 T^2$ at low T (below 60 K), with A_2 constant and A_1 decreasing with increasing x (and decreasing T_c (ref. 10)). This is the main reason why the coupling strength parameter $g(x)$ was assigned uniquely to the T -linear term in equation (1). Consistent with previous reports^{45,54}, $\rho(0, T)$ develops a stronger T dependence at higher T (>70 K). In-plane MR measurements were carried out in a 35-T Bitter magnet at HFML-FELIX in Nijmegen with the magnetic field applied perpendicular or parallel to the CuO₂ planes. The sample stage was glued with GE varnish to maintain its orientation.

Longitudinal MR

In a quasi-two-dimensional (2D) metal hosting Landau quasiparticles, a magnetic field oriented parallel to the direction of the injected current will generate a negligible Lorentz force, and as a result, the orbital MR is expected to be close to zero. Extended Data Fig. 2 compares the transverse and longitudinal MR for three overdoped concentrations— $x = 0.159$, 0.174 and 0.188 —at $T = 75$ K. Clearly, the longitudinal MR is negligible at this temperature, implying that the MR of LCCO is entirely orbitally driven (at lower temperatures, a non-zero MR appears that we attribute to the emergence of superconducting fluctuations). This anisotropy in the MR in electron-doped cuprates has been observed previously^{22,49,54,55}, in addition to an isotropic negative component in the longitudinal MR in PCCO for $x \leq x_{\text{FSR}}$ that was attributed to AFM order⁵⁵. In the p-doped cuprates Bi₂Sr₂CuO_{6+ δ} and Tl₂Ba₂CuO_{6+ δ} , by contrast, the longitudinal MR is found to have a similar magnitude as the transverse MR⁴⁸, behaviour that is inconsistent with conventional (Boltzmann) transport theory.

Boltzmann formalism in electron-doped cuprates

The purely orbital nature of the MR in overdoped LCCO motivates us to examine its T and H dependencies within a standard Boltzmann treatment, although modified to take into account the presence of a marked in-plane anisotropy of the scattering rate. We use the SCTIF applied to a 2D metal in the time basis to avoid divergences that can occur as $B \rightarrow 0$. The SCTIF is the general solution of the Boltzmann transport theory equation within the relaxation time approximation and is used to determine the response of quasiparticles to an external electromagnetic field at arbitrary field strengths. The expression for conductivity within the time basis, after neglecting the temperature broadening of the FS, is^{56,57}

$$\sigma_{ij} = \frac{1}{4\pi^3} \int_{\text{FS}} d^2k \frac{1}{\hbar v_F} q v_i \int_0^\infty q v_j (-t) P_\phi(t) dt, \quad (2)$$

where v_{ij} are the velocity components with $\{i, j\} = \{x, y\}$. $P_\phi(t)$ is the probability of a quasiparticle starting at angle ϕ to survive for time t , defined as

$$P_\phi(t) := \exp \left[- \int_0^t \frac{dt'}{\tau(t')} \right], \quad (3)$$

where $\tau(t)$ is the velocity relaxation time, and all other parameters take their usual values. The role of the magnetic field can be made explicit within the SCTIF formalism by making a change of variables from t to φ , the (azimuthal) angle swept around the FS under the influence of the field $\mathbf{H} \parallel \mathbf{c}$. Writing $\varphi = \omega_c t$, where $\omega_c(\varphi) = eBk_F(\varphi) \cdot \vec{v}_F(\varphi) / \hbar k_F^2(\varphi)$, we can re-express the SCTIF formalism as

$$\sigma_{ij} = \frac{1}{4\pi^3} \int_{\text{FS}} d^2k \frac{1}{\hbar v_F} q v_i \int_0^{2\pi} d\varphi \frac{q v_j(\varphi)}{\omega_c(\varphi)} \int_0^\infty d\varphi' \exp \left(\int_\varphi^{\varphi'} d\varphi'' \frac{-\varphi''}{\omega_c(\varphi'') \tau(\varphi'')} \right) \quad (4)$$

The key parameter that determines the H dependence of the MR is the mean free path $\ell(v_F \tau)$, particularly its anisotropy. The FS of LCCO does not exhibit sharp corners and is not in close proximity to a van Hove singularity. Thus, the k dependence of v_F (which incorporates the variation in the density of states around the FS) is anticipated to be limited. Indeed, for $x = 0.188$, where the FS curvature is most pronounced, v_F varies by $\leq 10\%$ (ref. 26). It is, thus, the anisotropy of τ that is expected to dominate the MR. A standard tensor inversion is then used to convert the magnetoconductivity to MR.

Given the strong magnetic order observed below x_{FSR} , spin correlations probably manifest at all x even when their obvious signatures in the MR vanish. Such spin correlations may necessitate going beyond the relaxation time approximation by, for example, including the effects of CVCs, as highlighted in the main text. These CVCs have been predicted to give rise to a non-Fermi-liquid T dependence of the resistivity and have a crucial impact on the Hall effect³⁰, as discussed later.

Tight-binding parameterization

The general tight-binding model for 2D systems with tetragonal symmetry is given by

$$E(k) = \varepsilon_0 - 2t(\cos(k_x a) + \cos(k_y a)) - 4t'(\cos(k_x a) \cos(k_y a)) - 2t''(\cos(2k_x a) + \cos(2k_y a)) \quad (5)$$

where t , t' and t'' are the nearest, next-nearest and next-next-nearest-neighbour hopping parameters, and ε_0 is a measure of the chemical potential. We use the conditions given in ref. 26 in which $t = 0.3$ and $t''/t' = -0.5$ remain constant at all x , and assume that the Luttinger count $x = \frac{2A_{\text{FS}}}{A_{\text{BZ}}} - 1$ is satisfied to fix ε_0 and t'/t . Here $A_{\text{FS(BZ)}}$ is the area of the FS (BZ), respectively. The TB parameters for LCCO are listed in Extended Data Table 1. LCCO is tetragonal with an in-plane lattice parameter $a = 4.01$ Å. The $\hbar \vec{v} = \nabla_k E$ relation is used to obtain the components of the velocity in equation (2), leaving τ as the only parameter that needs to be explicitly defined.

Simulations

Although the form of the scattering rate $\tau^{-1}(\varphi, T, x)$ (redefined as $\Gamma(\varphi, T, x)$ hereafter) expressed in equation (1) is a result of extensive modelling, it contains only those components required to capture all the key features of the experimental resistivity and MR data (outside of the FSR regime) as a function of field strength, temperature and doping with a single doping-dependent variable parameter $g(x)$. In this section, we justify the inclusion of each component in $\Gamma(\varphi, T, x)$ and present alternative scenarios that fail to capture all of the key features of the MR.

As discussed in the main text, there are two primary candidates for the dominant, anisotropic inelastic scattering rate in overdoped LCCO: spin fluctuations and short-range charge correlations. According to a recent angle-resolved photoemission spectroscopy (ARPES) study, the nodal regions of the FS in LCCO exhibit non-Fermi-liquid behaviour⁵⁸, probably due to their close proximity to the AFM zone boundary. There are no known wavevectors connecting two nodal points, however, and no experimental evidence so far for hotspot scattering at the nodes. Hence, we do not consider this scenario here.

Within the spin fluctuation scenario, the core element of $\Gamma(\varphi, T, x)$ is a four-fold symmetric, double-peaked component that gives rise to eight broad peaks, each located at the intersection of the underlying FS and the AFM BZ boundary (Fig. 3) and reproduced in Extended Data Fig. 3b1 (note that the loci of these hotspots shift in accordance with the doping level). For simplicity, both the T -linear and T^2 components of $\Gamma(\varphi, T, x)$ are assigned this anisotropic form, though only the T -linear term incorporates $g(x)$, as motivated by the derivatives of the zero-field resistivity presented in Extended Data Fig. 1c. Combined with an isotropic impurity scattering rate Γ_0 of appropriate magnitude, this T -dependent scattering rate (labelled hereafter Γ_T) captures the most striking feature of the MR data, namely, the growth of $\Delta\rho/\rho(0)$ with increasing T (Extended Data Fig. 3a3). Recall that Kohler's rule would predict an opposite T dependence for $\Delta\rho/\rho(0)$. The H dependence of the MR is also reasonably well captured by $\Gamma_0 + \Gamma_T$, except two important details. First, at low fields, the derivatives fan out from $H = 0$ (Extended Data Fig. 3b4), whereas in the data (Extended Data Fig. 3a4), the initial slopes are essentially T independent. Second, the field dependence tends towards saturation as the anisotropy in Γ_T becomes washed out by ever-increasing cyclotron motion at low T , and at $T = 50$ K, the MR is purely quadratic. In reality, the MR of overdoped LCCO remains H linear up to fields as high as 65 T (ref. 24). This observation justifies the inclusion in our simulation of loci ('hotspots') on the FS that impede cyclotron motion¹⁹ at all temperatures in this range.

Extended Data Fig. 3c,d considers a combination of spin and charge fluctuation contributions to $\Gamma(\varphi, T, x)$ to illustrate the fact that adding extra complexity does not necessarily improve the simulations. In Extended Data Fig. 3c, for example, we add small hotspots at the antinodal regions to mimic additional scattering off charge fluctuations between the antinodes (Fig. 3, grey double-headed arrows)^{50,51}.

The resulting simulations show an MR that is quadratic at low fields but which rapidly transitions towards saturation, as shown in the 20K MR curve (Extended Data Fig. 3c3) and its derivative (Extended Data Fig. 3c4). This tendency is due to the fact that the SCTIF concerns the average scattering rate around the FS; although this scenario contains substantive and complex anisotropy, such a combination of two anisotropic components generates a total scattering rate that is, in fact, on a coarse level, more isotropic, thereby explaining the tendency towards MR saturation.

In Extended Data Fig. 3d, we present the converse situation with dominant scattering from charge fluctuations, combined with subdominant hotspots at the AFM BZ boundary. In this case, the anisotropic inelastic component is a broad peak centred around the antinodes. This simulation matches more closely with the experimental curves; the initial slopes are essentially independent of temperature in this scenario (Extended Data Fig. 3d4), although once again, the simulated MR shows a tendency towards saturation at higher field strengths, particularly at lower T .

Finally, to convince the reader that hotspots alone are insufficient to reproduce the experimental data, Extended Data Fig. 3e shows the results of a simulation in which the anisotropy in $\Gamma(\varphi)$ is contained solely in the sharp hotspots positioned at the AFM zone boundary. Although the T dependence of $\rho(0, T)$ is reasonably well-matched in this scenario, (Extended Data Fig. 3e2), the T dependence of the MR is inverted (Extended Data Fig. 3e3). This demonstrates the clear need for the inclusion of Γ_T in the full expression for $\Gamma(\varphi, T, x)$.

The scattering rate parameters that provide the best overall fits to the $x = 0.159$ data based on a pure spin or charge fluctuation scenario (Fig. 4) are listed in Extended Data Table 2. The hotspot scattering rate τ_{HS}^{-1} is constructed as a narrow Gaussian peak whose height and width are also given. These hotspot peaks have a similar magnitude as the inelastic scattering rate, resulting in an overall Γ that is either peaked at the AFM zone boundary or at the antinodes, both scenarios providing the anisotropy that is required to reproduce the MR at low T . This parameterization is also consistent with the magnitude of the residual resistivity. Since τ_{HS}^{-1} peaks sharply around the intersections of the FS and the AFM BZ, the dominant contribution to the residual resistivity is τ_{imp}^{-1} . Ignoring the slight in-plane anisotropy in k_F and v_F , we can rewrite $\rho_0 = (2\pi\hbar c/e^2 k_F v_F) \tau_{\text{imp}}^{-1}$. Inserting the values for k_F and v_F from Fig. 4a(i) and for τ_{imp}^{-1} from Extended Data Table 2, we arrive at $\rho_0 \approx 15 \mu\Omega \text{ cm}$. Using the hotspot parameters quoted in Extended Data Table 2, we estimate that the angle-averaged $\tau_{\text{HS}}^{-1} \approx 4 \text{ ps}^{-1}$, giving a total $\rho_0 \approx 17.2 \mu\Omega \text{ cm}$. From Extended Data Fig. 1, we estimate that the actual value of $\rho_0 \approx 17.5 \mu\Omega \text{ cm}$.

Incorporation of CVCs into the Boltzmann formalism

In a 2D metal, the Hall conductivity σ_{xy} is determined by the 'Stokes' area swept out by the $\vec{\ell}(\varphi)$ vector as φ varies around the FS²⁹. The FS of LCCO (Figs. 3 and 4) has only hole-like curvature, meaning that $\vec{\ell}_F$ has the same circulation at all φ . In this circumstance, even the inclusion of strong (orders-of-magnitude) anisotropy in $\tau^{-1}(\varphi)$ cannot cause σ_{xy} , and therefore R_H , to become negative. In LCCO, for $x > x_{\text{FSR}}$, however, $R_H(T)$ drops monotonically from close to its isotropic (Drude) value towards zero and even becomes negative at elevated temperatures^{10,27} (Extended Data Fig. 4). Although a two-band model is able to account for the behaviour of the Hall resistance, it neither respects the FS measured in ARPES nor can it simultaneously account for the MR⁵⁹. In the absence of any additional FSR, for which there is no reported evidence, the only effective way of explaining the sign change in $R_H(T)$ within a single fluid picture is a breakdown of the relaxation time approximation and the incorporation of CVCs into the calculation for σ_{xy} (ref. 28).

In this section, we develop a semiquantitative picture of the evolution of $R_H(T)$ by incorporating the effect of CVCs into the analysis of the Hall response. For these simulations, we use the Jones–Zener expansion of the Boltzmann transport equation. Although this form of the Boltzmann equation is only applicable at low magnetic fields, R_H itself has only been measured so far at relatively low field strengths. Moreover, this form allows the use of analytical expressions for most of the key quantities. We start by introducing analytical expressions for $k_F(\varphi)$ and $v_F(\varphi)$ that simplify the TB parameterizations above and preserving their four-fold symmetry:

$$k_F(\varphi) = k_{F0}(1 + \alpha \sin^2(2\varphi)) \quad (6)$$

$$v_F(\varphi) = v_{F0}(1 + \beta \sin^2(2\varphi)) \quad (7)$$

Here $k_{F0} = 5.424 \times 10^9 \text{ m}^{-1}$, $\alpha = 0.11$, $v_{F0} = 4.946 \times 10^5 \text{ m s}^{-1}$ and $\beta = 0.03$. Note that in our simulations of the Hall response, the variation in $v_F(\varphi)$ at the antinodes has been flattened slightly relative to equation (7) to approximate the TB parameterization more closely. This has no impact on $\rho(H, T)$. $v_x(\varphi)$ and $v_y(\varphi)$ are then obtained from $v_F(\varphi)$ using the following analytical expressions:

$$\begin{aligned} v_x(\varphi) &= v_F(\varphi) \cos(\varphi - \gamma) \\ v_y(\varphi) &= v_F(\varphi) \sin(\varphi - \gamma) \end{aligned} \quad (8)$$

where⁶⁰

$$\gamma(\varphi) = \tan^{-1}(\partial [\ln[k_F(\varphi)]] / \partial \varphi) \quad (9)$$

In electron-doped cuprates close to commensurate AFM order, the dominant scattering vector is $\vec{Q} = (\pi, \pi)$. When CVCs are introduced,

momenta linked through $\vec{Q}$ are heavily mixed through strong scattering, as a result of which the total current J_k becomes approximately parallel to $v_k + v_{k+Q}$ and is, thus, no longer perpendicular to the FS. Previously, Kontani used the linear response theory to derive the following expression to describe how $\vec{J}_k$ transforms in proximity to the AFM hotspot³⁰:

$$\vec{J}_k = \frac{1}{1 - \epsilon_k^2} (\vec{v}_k + \epsilon_k \vec{v}_{k+Q}) \quad (10)$$

where $\epsilon_k \approx (1 - c\xi_{\text{AFM}}^2) < 1$, ξ_{AFM} is the AFM correlation length and c is a constant. To incorporate CVCs into our Boltzmann framework, we first substitute this expression for $\vec{J}_k(\vec{v}_k)$ into the kinetic equation and recalculate v_x and v_y :

$$\begin{aligned} v_x^*(\varphi) &= \frac{1}{1 - \epsilon(\varphi)^2} (v_x(\varphi) + \epsilon(\varphi)v_x(\varphi')), \\ v_y^*(\varphi) &= \frac{1}{1 - \epsilon(\varphi)^2} (v_y(\varphi) + \epsilon(\varphi)v_y(\varphi')), \end{aligned} \quad (11)$$

where $\varphi' = \varphi - (\pi, \pi)$. This process effectively links the hotspots through the AFM Q vector. For each value of x , the angles φ_n associated with the hotspots are determined by the locations at which the FS crosses the AFM zone boundary (for example, Extended Data Fig. 5c,d). For $\epsilon(\varphi)$, we assume a Gaussian with a fixed breadth and height (but, for simplicity, ignore the prefactor $1/(1 - \epsilon^2)$), using the following expression:

$$\epsilon(\varphi) = \exp \frac{-(\varphi - \varphi_n)^2}{2\sigma^2} \quad (12)$$

with $\sigma = 0.07$ (chosen to approximate the anisotropy expressed in equation (1)). We then substitute these expressions for v_x^* and v_y^* into the Jones–Zener expansions for σ_{xx} and σ_{xy} :

$$\sigma_{xx} = \frac{e^2}{2\pi^2\hbar c} \int_0^{2\pi} \frac{k_F(\varphi) v_x^*(\varphi) \cos(\varphi - \gamma)}{\Gamma(\varphi) \cos \gamma} d\varphi \quad (13)$$

$$\sigma_{xy} = \frac{-e^3}{2\pi^2\hbar^2 c} \int_0^{2\pi} \frac{v_x^*(\varphi)}{\Gamma(\varphi)} \frac{\partial [v_y^*(\varphi)/\Gamma(\varphi)]}{\partial \varphi} d\varphi \quad (14)$$

where $c = 6.2 \text{ \AA}$ is the c -axis lattice spacing. Finally, for $\Gamma(\varphi, T, x)$, we adopt the same form as in equation (1), which is reproduced here combined with the (isotropic) impurity scattering rate Γ_0 :

$$\Gamma(\varphi, T, x) = \Gamma_0 + \Gamma_{\text{HS}}(\varphi) + (g(x)\alpha_1 T + \alpha_2 T^2) \sin^2 \left[2 \left(\varphi \pm \left(\frac{\pi}{4} - \text{HS} \right) \right) \right] \quad (15)$$

where Γ_{HS} is the additional sharp Gaussian contribution at each hotspot whose parameters are given in Extended Data Table 2.

Anisotropic scattering term in LCCO

Extended Data Fig. 5a,b shows $v_x(\varphi)$ and $v_y(\varphi)$, respectively, together with their CVC-modified counterparts $v_x^*(\varphi)$ and $v_y^*(\varphi)$ for the parameterizations defined in equation (11). The expression for the Hall conductivity σ_{xy} of a 2D metal²⁹ (equation (14) and reproduced in the box at the top of the figure) is set by the product of $\ell_x(\varphi)$ and $\partial \ell_y(\varphi)/\partial \varphi$. In the absence of anisotropy in $\ell(\varphi)$ or any CVCs, $\ell_x \times \partial \ell_y$ is always positive. Introducing the CVCs, however, creates a window in φ , where $v_x^*(\varphi)$ and $\partial v_y^*/\partial \varphi$ have opposite sign, as indicated by the vertical shaded bars in both panels. Within the Ong construction, this sign change arises due to a change in circulation of the $\vec{v}$ vector around the FS. These regions of opposite circulation are highlighted in Extended Data Fig. 5c,d (thick solid arrows). Correspondingly, these regions can contribute a component to σ_{xy} of opposite sign, even when the scattering time τ is isotropic. The width of these regions is set by the width of the ‘influence’

of the CVCs. Incorporating anisotropy into $\tau(\varphi)$ (or equivalently $\Gamma(\varphi)$) will act to enhance or reduce this effect, depending on the form of $\Gamma(\varphi)$ itself.

For $R_H(T)$ to drop in magnitude with increasing T relative to its isotropic, low- T value, we require that ℓ in those regions in which the circulation of $\ell(\varphi)$ is opposite to that of φ to become longer—a scenario we label Γ_N in Extended Data Fig. 5c. Conversely, if ℓ is shorter at the antinodes, the opposite effect occurs and σ_{xy} (and hence R_H) will become larger as the anisotropy in the total scattering rate grows. This scenario, labelled Γ_{AN} , is shown in Extended Data Fig. 5d. The former situation (Γ_N) is relevant for the case in which spin fluctuation scattering is dominant because the nodes are notably closer to the AFM hotspots. $\Gamma_{AN}(\varphi)$, on the other hand (Extended Data Fig. 5d), is more likely to be associated with small- $\vec{Q}$ charge fluctuations or strong electron–phonon scattering, as suggested by recent ARPES measurements on LCCO that report the presence of a ubiquitous antinodal kink at a characteristic (doping-independent) energy scale ($\sim 45 \text{ meV}$). Its ubiquity points towards coupling to phonons⁶¹. Although this phonon-derived mode may be manifest in the single-particle response, its impact on the transport properties has not yet been verified.

An inspection of Extended Data Fig. 5c,d provides a qualitative understanding of which form of $\Gamma(\varphi)$ will act to reduce R_H with increasing T . As explained in the figure caption, for $\Gamma_N(\varphi)$, the counter-circulation of $\vec{v}(\varphi)$ occurs in the regions in which ℓ is the longest. Hence, as anisotropy in $\Gamma_N(\varphi)$ grows with increasing T , the negative contributions to σ_{xy} also grow, causing $R_H(T)$ to drop. Conversely, were Γ_{AN} to be the appropriate form of the T -dependent anisotropic component, $R_H(T)$ would increase with increasing T .

This analysis is put on a more quantitative footing in Extended Data Fig. 6, where R_H is calculated within the same Jones–Zener formalism for different forms of $\Gamma(\varphi)$. For all panels, we use a parameterization for $k_F(\varphi)$ and $v_F(\varphi)$ that is consistent with those used for $x = 0.159$ in Fig. 4. Extended Data Fig. 6a shows the isotropic case: $\Gamma(\varphi) = \Gamma_0$. In this instance, we obtain $R_H = +8.9 \times 10^{-10} \text{ m}^3 \text{ C}^{-1}$, about 20% higher than the Drude estimate for a (hole) carrier density of $1 - 0.159 = 0.841$ per unit cell and in excellent agreement with the data shown in Extended Data Fig. 4a for $x = 0.16$ at low T , even with the inclusion of CVCs. For $\Gamma(\varphi) = \Gamma_0 + \Gamma_{\text{HS}}$ (Extended Data Fig. 6b), we find $R_H = +5.6 \times 10^{-10} \text{ m}^3 \text{ C}^{-1}$, showing that the addition of hotspot scattering leads to a slight reduction in R_H .

As argued above, a larger R_H is found when the scattering rate (between the hotspots) is larger in the antinodes. Extended Data Fig. 6c depicts a form of $\Gamma(\varphi) = \Gamma_0 + \Gamma_{\text{HS}} + \Gamma_{AN}$. This form of scattering could reflect a scenario in which strong charge fluctuations coexist with the AFM hotspots. In this case, $R_H = +1.03 \times 10^{-9} \text{ m}^3 \text{ C}^{-1}$, confirming that additional scattering at the antinodes leads to a further increase in R_H in the presence of CVCs. Hence, as anisotropy in $\Gamma(\varphi)$ grows with increasing T , so will R_H . This is opposite to what is observed experimentally^{10,27}. Clearly, to simulate a vanishing R_H , the anisotropic component of $\Gamma(\varphi)$ needs to be inverted.

This is done in Extended Data Fig. 6d, where now $\Gamma(\varphi) = \Gamma_0 + \Gamma_{\text{HS}} + \Gamma_N$. Using the same parameters that were used to model the MR in the main text (in the absence of CVCs), we find $R_H = +4.3 \times 10^{-10} \text{ m}^3 \text{ C}^{-1}$, which is about half the value found for Γ_0 only.

Finally, Extended Data Fig. 6e shows that removing the AFM hotspots and considering only CVCs and anisotropic scattering created by short- Q charge fluctuations leads to a marked enhancement in R_H . This is due to the fact that at the loci on the FS where the mixing of states occurs (that is, where the CVCs are finite), (i) the mean free path is strongly suppressed and (ii) the FS is essentially flat with very little curvature. Indeed, despite extensive analysis, we were unable to find a simulation based purely on dominant short- Q charge fluctuation scattering that could reproduce both the growth in the MR and the concomitant drop in $R_H(T)$ with increasing temperature. We, thus, do not consider this scenario further in this article. Finally, in

Supplementary Fig. 5, we show that the incorporation of modified velocities (from spin fluctuation scattering) into the SCTIF has a negligible effect on the modelling of the MR.

In summary, we find that in addition to CVCs due to AFM spin fluctuations, the inclusion of an associated Γ_N with a magnitude that leads to a doubling of the MR (Fig. 4) also causes R_H to halve. This form of $\Gamma(\varphi, T, x)$, thus, enables the two key T dependencies in the magnetotransport of LCCO beyond $x > x_{\text{FSR}}$ to be simulated within the same parameterization. This result also echoes an earlier study of the frequency (ω) dependence of the Hall angle in overdoped PCCO where a sign change of $R_H(T, \omega)$ in an FS with only a singular curvature was also interpreted as evidence for AFM CVCs²⁸. This analysis also suggests that it should be possible to reproduce both $R_H(T, x)$ and the MR from the ARPES-derived FS of LCCO using a single set of parameters. Given the large number of variables (width, height of hotspot, precise form of the anisotropic component and so on), however, here we simply present this semiquantitative analysis to justify the phenomenological form of $\Gamma(\varphi)$ and close this section with some preliminary thoughts on its possible physical origin.

The phenomenological form of $\Gamma(\varphi)$ —necessitated primarily by the unusual growth in $\Delta\rho/\rho(0)$ in LCCO over a temperature range in which ℓ itself is diminishing—has three principal features: T linearity; large, smooth anisotropy; and maxima at the AFM BZ boundary. The general approach to account for such a combination of features is to invoke the presence of low-lying bosonic excitations. The most natural candidates for these excitations are magnon modes in the vicinity of $\vec{Q} = (\pi, \pi)$. As shown in Fig. 3, the node is three times closer to the AFM zone boundary than the antinode. Thus, any spectral broadening of the (π, π) mode (due to magnon lifetime effects) or any thermal occupation of modes up to about 10 meV at 50 K will disproportionately widen the hotspot towards the node. Both effects have been observed in electron-doped cuprates^{46,61,62} and, when combined with the analysis presented above, offer a viable explanation as to why $R_H(T)$ decreases with increasing T but the MR increases (at least up to 70 K). This also may explain the T linearity of $\Gamma(\varphi, T, x)$.

A final key question to address, though, is how scattering off these magnon modes can induce such a large anisotropy in Γ . Indeed, recent calculations based on the dynamical cluster approximation applied to overdoped cuprates show that scattering off low-lying spin excitations generates only mild anisotropy in the scattering rate⁶³. One possible source for the enhancement seen in the experiment is small-angle scattering. Cuprates are known to exhibit primarily out-of-plane disorder in the charge-transfer layers rather than in CuO_2 planes, resulting in small-angle scattering on the electrons^{64,65}. Small-angle scattering can be understood as Brownian motion or as an effective diffusion process across the FS, and the expected time taken to diffuse into the nearby hotspots scales as the sector size squared (where the hotspots act as ‘sinks’ similar to the necks in elemental Cu)^{66,67}. If the hotspot in LCCO indeed widens predominantly towards the node with increasing temperature, then the shrinking nodal sector can cause a marked increase in the efficacy of small-angle scattering to relax the velocity, thereby enhancing the anisotropy between the nodal and antinodal regions. Thus, one possible interpretation of the increasing anisotropy in LCCO with temperature through the $g(x)$ term is that small-angle scattering exacerbates the anisotropic scattering. Although modelling this scenario is beyond the present work, we suggest that hotspot broadening combined with small-angle scattering may explain the origin of the unusual T -dependent anisotropy in Γ obtained here.

Before closing, we highlight two aspects of the magnetotransport in LCCO that are not yet understood. First among these is the T^2 resistivity that extends to high T (refs. 45, 59). As shown in Fig. 1, this behaviour contrasts markedly with the T linearity found in p-type cuprates at high T . The fact that $\rho_{ab}(T)$ in LCCO extends beyond the estimated Mott–Ioffe–Regel limit suggests that this superlinear behaviour may arise from a gradual loss of carriers or a transfer of low-energy spectral

weight with increasing T . The latter could be revealed, for example, in optical conductivity measurements at elevated temperatures⁶⁶.

The other outstanding issue is the persistence of H -linear MR to the lowest T (ref. 24). Within the model presented above, the H -linear MR is a result of impeded cyclotron motion due to the presence of hotspots. Of course, there may be a physical reason for this, for example, non-local (k -dependent) impurity potential scattering or the presence of a quantum critical phase and associated residual scattering off quantum fluctuations associated with the AFM order, but for the latter, it is not immediately obvious why it should persist down to low T . An alternative explanation is that this residual, small H -linear MR arises due to the compositional variation along the film created by the combinatorial synthesis itself. Longitudinal variation in the Hall voltage is a long-standing explanation for H -linear MR in metals⁶⁸ and, indeed, has recently been proposed in the context of cuprates⁶⁹. We note, however, that H -linear MR is also observed at low T in LCCO films that are not grown combinatorially²⁴. With this in mind, it is tempting to attribute the persistence of a sizable $1/\tau_{\text{HS}}$ down to low temperatures to the sublinear behaviour of $1/\tau(\propto \sqrt{T})$ predicted near the hotspots⁷⁰ within spin fluctuation theory, though this is clearly an outstanding issue that awaits explanation based on future microscopic theories.

Evolution of the high-field slope over the full (T, x) range

As a further, final hint to the origin of the anisotropic scattering rate in LCCO, Extended Data Fig. 7 shows the high-field slope γ_1 , defined as the value of $d\rho/d\mu_0 H$ at 33 T, with respect to T and x . At all doping levels, irrespective of the AFM order, the slope of MR falls as $T < 75$ K. Within the AFM regime, the slope becomes negative below $T \approx 20$ K without discontinuities, suggesting that the origin of its T dependence across the entire T range is closely linked to the evolution of AFM correlations, given the negative slope is associated with the AFM order⁸. When $x > x_{\text{FSR}}$, the slope is positive at all T , but its magnitude also falls below 75 K before levelling off at 20 K.

The smooth variation in the MR and the distinct similarity across the AFM boundary strongly links the unusual T dependence of the MR below 75 K to AFM spin correlations that persist across the doping series. Note that the T dependence diminishes as x increases, reflecting the weakening spin fluctuations implied by the decrease in the scaling parameter $g(x)$. Alternatively, the monotonic x dependence is also evident from the gradual variation in $d\rho/d\mu_0 H$ (plotted versus H/T in Fig. 2h–i) with increasing x . The low- T negative component seen at low x also evolves continuously, becoming positive with no discontinuity near x_{FSR} .

Data availability

The data that support the findings of this study are available via a Radboud University dataset at <https://doi.org/10.34973/x0sx-6q62> (ref. 71).

Code availability

The code that supports the findings of this study is available via a Radboud University dataset at <https://doi.org/10.34973/x0sx-6q62> (ref. 71).

References

- Bach, P. L., Saha, S. R., Kirshenbaum, K., Paglione, J. & Greene, R. L. High-temperature resistivity in the iron pnictides and the electron-doped cuprates. *Phys. Rev. B* **83**, 212506 (2011).
- Jin, K., Zhang, X. H., Bach, P. & Greene, R. L. Evidence for antiferromagnetic order in $\text{La}_{2-x}\text{Ce}_x\text{CuO}_4$ from angular magnetoresistance measurements. *Phys. Rev. B* **80**, 012501 (2009).
- Shockley, W. Effect of magnetic fields on conduction—‘tube integrals’. *Phys. Rev.* **79**, 191 (1950).
- Chambers, R. G. The kinetic formulation of conduction problems. *Proc. Phys. Soc. A* **65**, 458 (1952).

58. Tang, C. Y. et al. Antinodal kink in the band dispersion of electron-doped cuprate $\text{La}_{2-x}\text{Ce}_x\text{CuO}_{4\pm\delta}$. *npj Quantum Mater.* **7**, 53 (2022).
59. Li, P., Balakirev, F. F. & Greene, R. L. High-field Hall resistivity and magnetoresistance of electron-doped $\text{Pr}_{2-x}\text{Ce}_x\text{CuO}_{4-\delta}$. *Phys. Rev. Lett.* **99**, 047003 (2007).
60. Hussey, N. E. The normal state scattering rate in high- T_c cuprates. *Eur. Phys. J. B* **31**, 495 (2003).
61. Wilson, S. D. et al. High-energy spin excitations in the electron-doped superconductor $\text{Pr}_{0.88}\text{LaCe}_{0.12}\text{CuO}_{4-\delta}$ with $T_c = 21\text{ K}$. *Phys. Rev. Lett.* **90**, 157001 (2006).
62. Fujita, M., Matsuda, M., Fak, B., Frost, C. D. & Yamada, K. Novel spin excitations in optimally electron-doped $\text{Pr}_{0.89}\text{LaCe}_{0.11}\text{CuO}_4$. *J. Phys. Soc. Jpn* **75**, 093704 (2006).
63. Bergeron, D., Hankevych, V., Kyung, B. & Tremblay, A.-M. S. Optical and d.c. conductivity of the two-dimensional Hubbard model in the pseudogap regime and across the antiferromagnetic quantum critical point including vertex corrections. *Phys. Rev. B* **84**, 085128 (2011).
64. Varma, C. M. & Abrahams, E. Effective Lorentz force due to small-angle impurity scattering: magnetotransport in high- T_c superconductors. *Phys. Rev. Lett.* **86**, 4652 (2001).
65. Abrahams, E. & Varma, C. M. Hall effect in the marginal Fermi liquid regime of high- T_c superconductors. *Phys. Rev. B* **68**, 094502 (2003).
66. Pippard, B. A. Longitudinal magnetoresistance. *Proc. R. Soc. Lond. A* **282**, 464 (1964).
67. Hussey, N. E., Takenaka, K. & Takagi, H. Universality of the Mott–Ioffe–Regel limit in metals. *Philos. Mag.* **84**, 2847 (2004).
68. Bruls, G. J. C. L. B. et al. Linear magnetoresistance caused by sample thickness variations. *Phys. Rev. Lett.* **46**, 533–536 (1981).
69. Singleton, J. Temperature scaling behaviour of the linear magnetoresistance observed in high-temperature superconductors. *Phys. Rev. Mater.* **4**, 061801 (2020).
70. Stokjovic, B. P. & Pines, D. Theory of the longitudinal and Hall conductivities of the cuprate superconductors. *Phys. Rev. B* **55**, 8576 (1997).
71. Duffy, C. M. et al. Evidence for spin-fluctuation-mediated superconductivity in electron-doped cuprates. Version 1. *Radboud University* <https://doi.org/10.34973/x0sx-6q62> (2025).

Acknowledgements

We acknowledge stimulating discussions with J. Ayres, G. Grissonnanche, A.-M. Tremblay, C. Iorio-Duval, L. Taillefer, A. Ataei and C. Proust during the course of preparing this manuscript. We

acknowledge support from HFML-FELIX, member of the European Magnetic Field Laboratory (EMFL).

Author contributions

S.J.T. and Q.H.C. grew the LCCO thin films. C.M.D., Q.H.C. and J.S.Z. performed the transport measurements. C.M.D. and N.E.H. analysed the experimental data. C.M.D. performed the SCTIF simulations. R.D.H.H. and A.C. provided the code for the simulations. C.M.D. and N.E.H. wrote the manuscript with input from all authors. N.E.H. and K.J. supervised the project.

Funding

This work was supported by the European Research Council (ERC) under the European Union's Horizon 2020 research and innovation programme (grant agreement number 835279-Catch-22; C.M.D., A.G., R.D.H.H. and N.E.H.) and the Engineering and Physical Sciences Research Council (UK) grant EP/V02986X/1 (N.E.H.). T.S. and R.L.G. acknowledge support from the NSF grant DMR 2002658 and the Maryland Quantum Materials Centre. S.J.T., Q.H.C., J.S.Z. and K.J. acknowledge support from the NSFC grant (12225412 and 12374141), the National Key R&D Program of China (2022YFA1403900 and 2022YFA1403000), the Beijing Natural Science Foundation (JQ26007) and the Chinese Academy of Sciences President's International Fellowship Initiative (2025PG0007).

Competing interests

The authors declare no competing interests.

Additional information

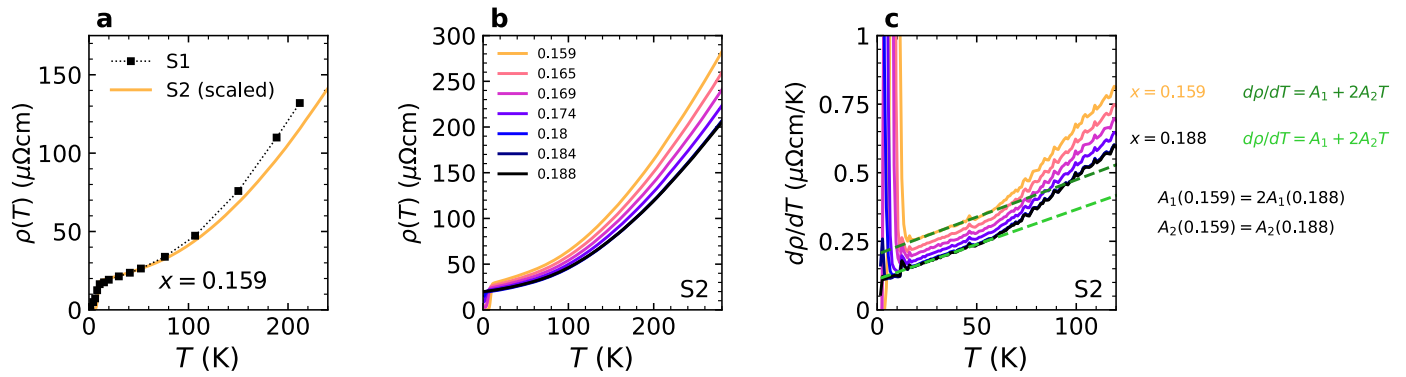
Extended data is available for this paper at <https://doi.org/10.1038/s41567-026-03454-6>.

Supplementary information The online version contains supplementary material available at <https://doi.org/10.1038/s41567-026-03454-6>.

Correspondence and requests for materials should be addressed to C. M. Duffy, Q. H. Chen, K. Jin or N. E. Hussey.

Peer review information *Nature Physics* thanks Hiroshi Kontani and the other, anonymous, reviewer(s) for their contribution to the peer review of this work. Peer reviewer reports are available.

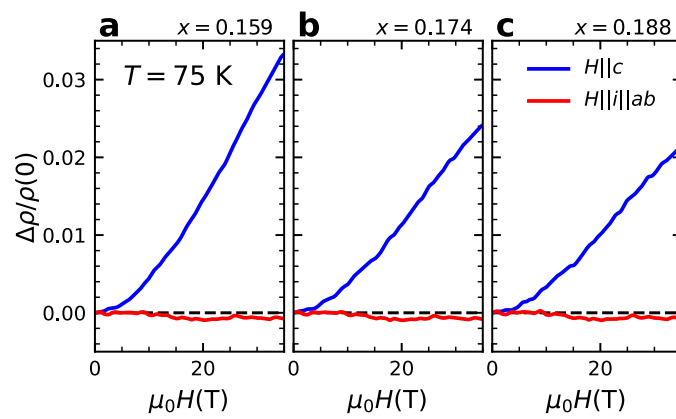
Reprints and permissions information is available at www.nature.com/reprints.



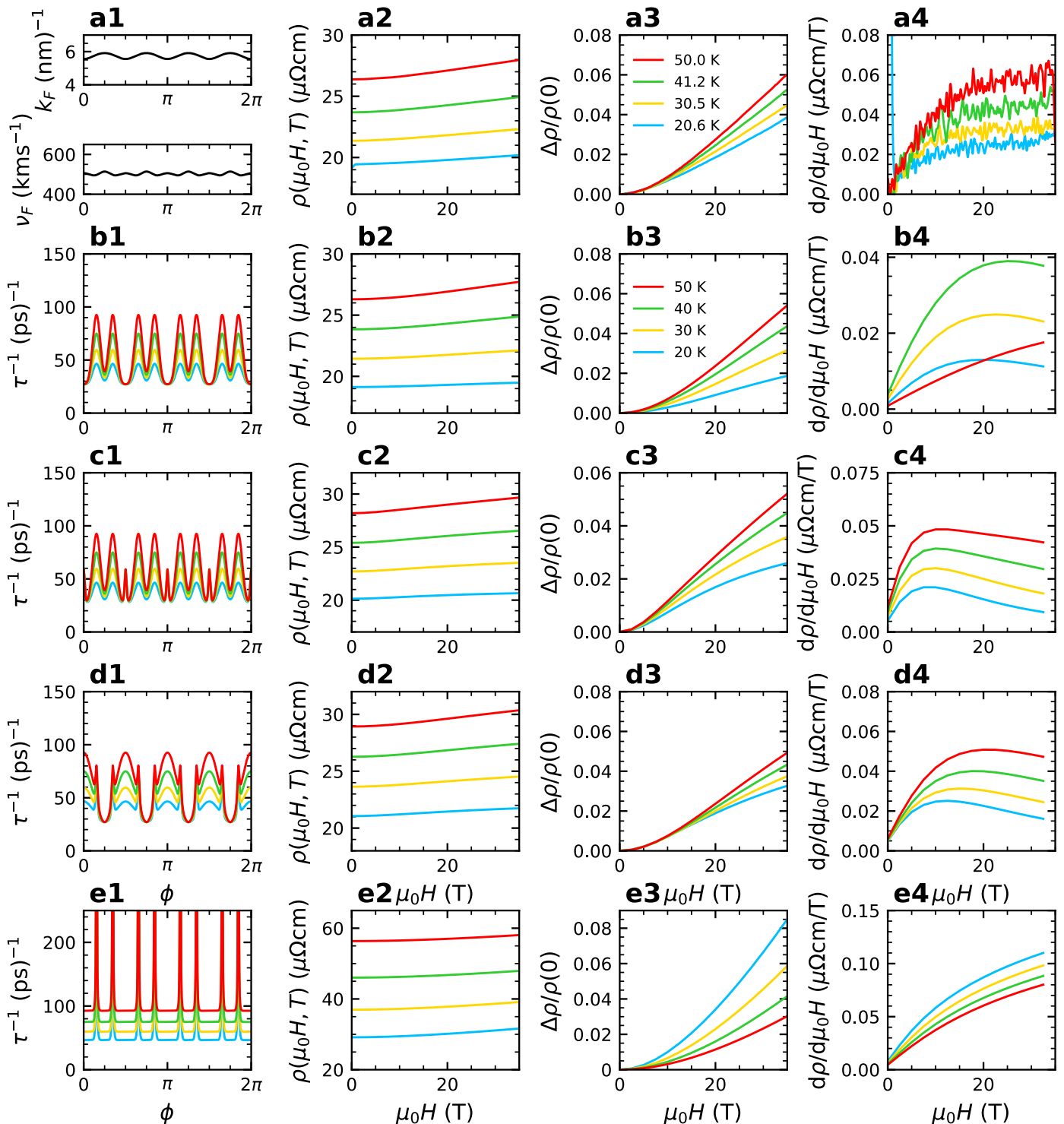
Extended Data Fig. 1 | Zero-field resistivity of overdoped LCCO films. (a)

$\rho(0, T)$ at $x = 0.159$ in the two films. Note that the resistivity values of S2 have been divided by 1.5 in order to match those of S1 below 70 K. **(b)** Representative $\rho(0, T)$ curves for the LCCO film S2 (unscaled). The form of $\rho(0, T)$ is the same across all channels. **(c)** Corresponding derivatives up to 125 K. The finite intercept as $T \rightarrow 0$

indicates a persistent T -linear component whose strength diminishes with x . Between 15 K and 60 K, the derivatives also reveal a constant T^2 component of $\rho(0, T)$ across the doping range. The discontinuous change in the derivative of the resistivity at $T \approx 60$ K (ref. 45) coincides with the temperature at which the MR and the value of the H -linear slope are maximal.

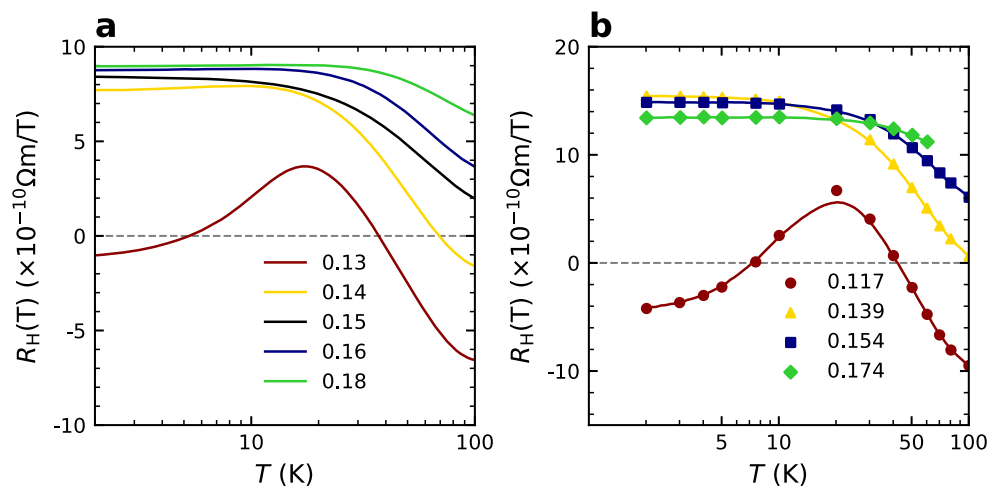


Extended Data Fig. 2 | Orientation dependence of the MR. The MR of LCCO with $H||c$ (blue) and $H||l||ab$ (red) for three channels (a) $x = 0.159$, (b) $x = 0.174$ and (c) $x = 0.188$ at $T = 75$ K. The MR is anisotropic with respect to field orientation and is close to zero in the parallel configuration as is expected from conventional theories of transport.



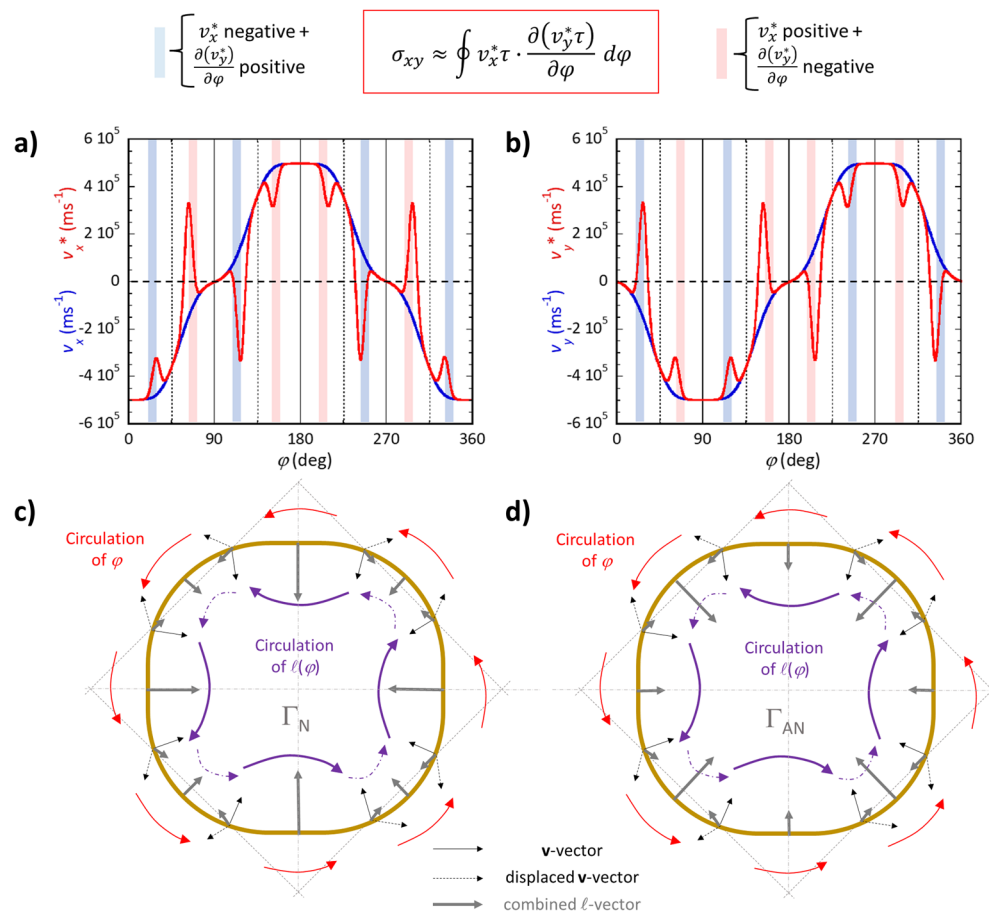
Extended Data Fig. 3 | (Inferior) simulations with alternative scattering rates for $x = 0.159$. Reproduction of row **a**) from Fig. 4 of the main article, for comparison. **(b)** Anisotropic scattering (peaked at the AFM Brillouin zone boundary) without hotspots. While the zero-field resistivity is in excellent agreement with the experimental data, from panel b3 it is clear that the T -dependence of the MR at high magnetic fields cannot simultaneously be captured without the inclusion of hotspots. Furthermore, from the derivatives in panel b4, the simulated data at 50 K is approximately quadratic in field at all field strengths, whereas the data show a strong linear component. **(c)** Same as **(b)** but with (small) hotspots positioned at the antinodes. In this case, the anisotropic

components are spread evenly across the FS. As a result, the anisotropy is washed out rapidly with increasing H and the MR tends towards saturation at all T . **(d)** Anisotropic scattering peaked at the antinodes combined with hotspots at the AFM zone boundary. Again, the MR has a stronger tendency to saturate at all T . **(e)** Isotropic inelastic scattering with stronger hotspot scattering (10^5 ps^{-1}). In this case, the T -dependence of the MR is inverted highlighting the need for the anisotropic component in $\Gamma(\phi, T, x)$. The stronger hotspot also amplifies the T -dependence of the zero-field resistivity in panel e2) beyond what is observed experimentally.



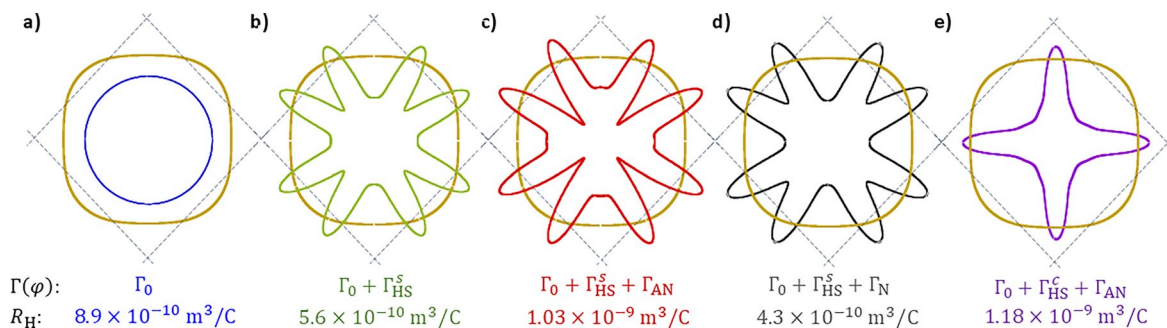
Extended Data Fig. 4 | Hall coefficient in LCCO. (a) $R_H(T)$ curves in non-combinatorial thin films of LCCO over the doping range $0.13 \leq x \leq 0.18$. Reproduced from ref. 27. (b) Selection of $R_H(T)$ curves in combinatorial LCCO

films over a similar doping range¹⁰. Note the propensity of $R_H(T)$ to decrease with increasing T and tend towards negative values even for overdoped samples beyond $x_{\text{FSR}} = 0.14$. Both sets of data were obtained in a field of 14 T.



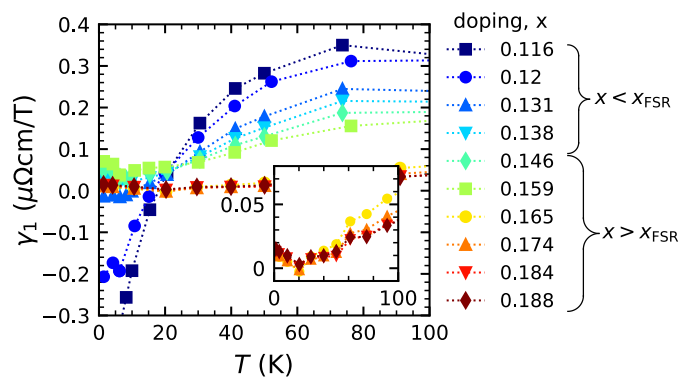
Extended Data Fig. 5 | Possible influence of current vertex corrections (CVCs) on the Hall response in LCCO. (a) Variation of $v_x(\varphi)$ (blue line) and its CVC-modified counterpart $v_x^*(\varphi)$ (red line). (b) Corresponding plot of $v_y(\varphi)$ and $v_y^*(\varphi)$. The shaded vertical bars represent those sections of the FS where the product of $v_x^*(\varphi)$ and $\partial v_y^* / \partial \varphi$ is negative, as described by the expressions above the panels. (c/d) Comparison of the impact on σ_{xy} of the change in circulation of $\ell(\varphi)$ due to CVCs for the case where (c) $\ell(\Gamma)$ at the nodes is shorter (greater) than at the anti-nodes (a scenario labelled Γ_N in panel c)) and (d) vice-versa (Γ_{AN}). The length of the grey arrows reflect the magnitude of $\ell(\varphi)$ at the symmetry points.

Wherever the circulation of $\ell(\varphi)$ opposes that of φ (that is anticlockwise), the contribution to σ_{xy} from that section is negative. σ_{xy} itself is dominated by those regions of the FS where the curvature in $\ell(\varphi)$ is largest. For LCCO, the largest curvature occurs around the zone diagonals. In panel c), the scattering rate at the nodes is greater than at the anti-nodes and so the (positive) contribution to σ_{xy} is reduced relative to the negative contribution from the more extended regions between the hotspots, leading to a reduction in σ_{xy} . Conversely, in d), σ_{xy} becomes larger (more positive) as the anisotropic scattering rate grows.



Extended Data Fig. 6 | Anisotropic scattering rates in LCCO. Candidate forms of $\Gamma(\varphi)$ – specified in the top line – that are inserted into the Jones-Zener expansion of the Boltzmann transport equation to calculate the corresponding R_H – specified in the second line below the plots. The forms of $\Gamma(\varphi)$ are: **a)** Γ_0 (isotropic); **b)** $\Gamma_0 + \Gamma_{\text{HS}}^s$ (including a set of peaks associated with the AFM spin hotspots on top of an isotropic component); **c)** $\Gamma_0 + \Gamma_{\text{HS}}^s + \Gamma_{\text{AN}}$ (in which the broad

anisotropic component is peaked at the anti-nodes); **d)** $\Gamma_0 + \Gamma_{\text{HS}}^s + \Gamma_{\text{N}}$ (in which the broad anisotropic component is peaked at along the zone diagonals); **e)** $\Gamma_0 + \Gamma_{\text{HS}}^c + \Gamma_{\text{AN}}$ (including a set of peaks associated with charge hotspots near the zone boundary). Only the combination of AFM hotspots plus a broad anisotropic component peaked along the zone diagonals (panel **d**) can generate a marked reduction in R_H .



Extended Data Fig. 7 | The link between the scattering rate anisotropy and AFM correlations. $\gamma_1 (= d\rho/d\mu_0 H/H = 33 \text{ T})$ vs. T for selected doping values across the range studied. For $x < x_{\text{FSR}}$, γ_1 falls smoothly below 75 K and becoming negative below 20 K without any discontinuities, indicating that the drop in the MR is a consequence of AFM order and that correlations remain present

at high temperatures. At higher x , γ_1 also falls below 75 K suggesting that AFM correlations still play a crucial role in the transport properties at elevated doping levels. The decrease in γ_1 with increased x is then a reflection of the weakening AFM order when doped further away from the AFM phase.

Extended Data Table 1 | Tight-binding parameters for LCCO based on the requirements listed in Tang et al.⁵⁷

x	ε_0	t	t'/t	t''/t'
0.1	-0.045	0.3	0.22	-0.5
0.125	-0.02059	0.3	0.225	-0.5
0.15	0	0.3	0.23	-0.5
0.175	0.033	0.3	0.235	-0.5
0.2	0.0615	0.3	0.24	-0.5

Using $t=0.3$ and $t''/t'=-0.5$, ε_0 and t'/t are found for $x=0.1-0.2$ through the Luttinger count.

Extended Data Table 2 | Scattering rate parameters used to fit $\rho(H,T)$ of sample S1 at $x=0.159$ in Fig. 4 of main article

τ_{imp}^{-1} (ps ⁻¹)	α_1 (ps ⁻¹ K ⁻¹)	α_2 (ps ⁻¹ K ⁻²)	HS peak (ps ⁻¹)	HS width
27	0.75	0.0112	50	0.05

τ_{imp}^{-1} ($= \Gamma_0$) is the isotropic impurity scattering rate, and α_1 (α_2) are the coefficients of the T -linear (T^2) components, respectively. The hotspots (in both the spin and charge fluctuation scenarios) are Gaussians with fixed height and width. Note that these parameters are kept fixed for all the other simulations presented in the main article.

Extended Data Table 3 | Doping dependence of $g(x)$ for S1 (top rows) and S2 (bottom rows)

x	$g(x)$
S1	
0.116	6.50
0.119	5.75
0.12	5.20
0.125	3.85
0.131	2.90
0.134	2.52
0.138	2.10
0.141	1.87
0.146	1.46
0.149	1.45
0.153	1.38
0.156	1.18
0.159	1.00
S2	
0.159	1.00
0.165	0.85
0.169	0.76
0.174	0.66
0.180	0.57
0.184	0.55
0.188	0.55

For S1, the anisotropic T -linear scaling factor $g(x)$ is obtained by normalizing to the value at $x=0.159$ (where the simulation fitting of the T - and H -dependence of the MR was carried out) and varying $g(x)$ such that it simulated the zero-field resistivity and MR magnitude at $T=40\text{ K}$ and $\mu_0 H=20\text{ T}$ of all other channels without changing any other parameter (see Fig. 5). Accordingly, $g(x)$ is found to be largest at the lowest doping and drops monotonically across the series. For S2, $g(x)$ is obtained from the coefficient of the T -linear resistivity shown in Extended Data Fig. 1c) and normalized such that the values of $g(x)$ in both films are the same at $x=0.159$. $g(x)$ is found to drop by roughly a factor of 2 across this doping range.