Saturation of resistivity and Kohler’s rule in Ni-doped La$_{1.85}$Sr$_{0.15}$CuO$_4$ cuprate

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We present the results of electrical transport measurements of La$_{1.85}$Sr$_{0.15}$Cu$_{1-y}$Ni$_y$O$_4$ thin single-crystal films at magnetic fields up to 9 T. Adding Ni impurity with strong Coulomb scattering potential to slightly underdoped cuprate makes the signs of resistivity saturation at $\rho_{sat}$ visible in the measurement temperature window up to 350 K. Employing the parallel-resistor formalism reveals that $\rho_{sat}$ is consistent with classical Ioffe-Regel-Mott limit and changes with carrier concentration $n$ as $\rho_{sat} \propto 1/\sqrt{n}$. Thermopower measurements show that Ni tends to localize mobile carriers, decreasing their effective concentration as $n \approx 0.15 - y$. The classical unmodified Kohler’s rule is fulfilled for magnetoresistance in the nonsupercconducting part of the phase diagram when applied to the ideal branch in the parallel-resistor model.

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Increasing evidence for well-defined quasiparticles in underdoped cuprates seems to corroborate a view that they are normal metals, only with small Fermi surface. Fermi-Dirac statistic underlying the quantum oscillations ($n$), single-parameter - quadratic in energy $\omega$ and temperature $T$ - scaling in optical conductivity $\sigma(\omega, T)$ (Ref. [2]), $T^2$ resistivity behavior extending over substantial $T$-region in clean systems [3] and fulfillment of typical for conventional metals Kohler’s rule in magnetotransport [4] are observations in favor of Fermi-liquid scenario.

On the other hand, in cuprates with significant disorder manifested by large residual resistivity $\rho_{res}$, pure $T^2$ resistivity dependence has not been reported so far as for Bi$_2$Sr$_2$CuO$_{6+\delta}$ [2, 3, 4] or observed only at relatively narrow doping- and $T$-region as in La$_{2-x}$Sr$_x$CuO$_4$ (LSCO) [5, 6]. The clear violation of Kohler’s scaling in underdoped LSCO and YBa$_2$Cu$_3$O$_{7-\delta}$ [4, 5] (although not necessarily meaning breakdown of Fermi-liquid description [10]) has served almost as a hallmark of their peculiar normal-state properties for two decades.

In contrast to overdoped cuprates where large cylindrical Fermi surface yields a carrier density $n=p+1$ ($p$ being doping level) [11, 13], the total volume of Fermi surface in underdoped systems is a small fraction of the first Brillouin zone and corresponds to $n=p$ through Luttinger’s theorem [14, 17]. This small $n$ should be reflected in zero-field transport. In normal metals, resistivity $\rho$ saturates in the vicinity of Ioffe-Regel-Mott limit $\rho_{IRM}$ where elastic mean free path $l_{min}$ becomes comparable to interatomic distance $a$ [18, 19]. In cuprates, however, signs of saturation are seen at $\rho_{sat}$ much larger than $\rho_{IRM}$ calculated from the semiclassical Boltzmann theory [20–23]. Moreover, $\rho(1000 \, \text{K})$ ($\sim \rho_{sat}$) in LSCO changes as $1/x$, while for $n \propto x$ (Ref. [24]) the theory predicts $\rho_{sat} \propto 1/\sqrt{x}$.

The above can be explained by breakdown of the quasiparticle picture due to strong inelastic scattering at high $T$ manifested by disappearing of a Drude peak in $\sigma(\omega)$ (Refs. [23, 25, 27]). In systems where impurity scattering dominates the carrier relaxation (quasiparticle decay) rate $1/\tau$, the Drude peak is centered at $\omega=0$ regardless of how strong scattering becomes [2, 23]. Electron-electron interactions make $\tau$ frequency-dependent but for Fermi-liquid-like $\omega^2$ dependence $\sigma(\omega)$ still peaks at $\omega=0$ (Ref. [2]). Thus large impurity-induced $\rho_{res}$ may facilitate approach to Ioffe-Regel-Mott limit in dc ($\omega=0$) LSCO transport at lower $T$ before the spectral weight is transferred to higher-energy excitations at larger $T$. Ni impurity is a good candidate because its strong Coulomb scattering potential in the CuO$_2$ planes allows to achieve large $\rho$ at moderately high $T$ [26, 27].

In this paper we report transport and thermopower measurements on slightly underdoped $x=0.15$ LSCO with added Ni impurity. The obtained $\rho_{sat}$ corresponds to the classical value for small Fermi surface and changes as $1/\sqrt{n}$. The Fermi-liquid quasiparticle picture holds in the nonsuperconducting part of the phase diagram, as revealed by $\rho \propto T^2$ dependence and classical Kohler’s rule for magnetoresistance, both hidden under the large resistivity of the system.

The 4-point transport measurements were carried out on the $c$-axis aligned single-crystal films grown on isostructural LaSrAlO$_4$ substrate by the laser ablation method from the polycrystalline targets [30, 31]. The thermopower, which is not sensitive to the grain boundaries and porosity effects in cuprates [32], was measured on the samples cut from the targets.

Figure 1 shows the systematic change in $\rho(T)$ with Ni, from superconducting $y=0$ specimen with midpoint $T_C=34.6 \, \text{K}$ to $y=0.08$ one exhibiting insulating behavior at low $T$. In high $T$, a change of slope in portion of $\rho(T)$ curves that increases with $T$ forers the approaching saturation. Variation of $d\rho/dT$, visible even in the $y=0$ data, authenticates the slope decreasing that becomes more pronounced with increasing $y$, as can be seen in Fig. 2(a). Resistivity in this region is described extremely well by the parallel-resistor formula

$$\frac{1}{\rho(T)} = \frac{1}{\rho_{id}(T)} + \frac{1}{\rho_{sat}} = \frac{1}{a_0 + a_1 T + a_2 T^2} + \frac{1}{\rho_{sat}}. \quad (1)$$
The ρsat term is the ideal resistivity in the absence of saturation 10 and the additive-in-conductivity formalism stems from existence of the minimal scattering time \( \tau_{\text{min}} \), equivalent to Ioffe-Regel-Mott limit, which causes the shunt ρsat to always influence ρ in normal metals 33-35. The formula was used for overdoped LSCO 35 but with the large-Fermi-surface ρsat value 34 as a fixed parameter. The excellent fits of Eq. 1 with all free parameters including ρsat to \( \rho(T) \) in 150 K-350 K interval are depicted in upper inset to Fig. 1. Extending the fit interval downwards to lower \( T \) for small \( y \) does not change significantly the obtained fitting parameters, which are presented in Fig. 2 (b-d).

The Ni-induced residual resistivity, calculated as \( 1/\rho_{\text{res}}=1/\rho_{\text{sat}}+1/a_{2} \), accelerates with \( y \) such that at large \( y \) substantially exceeds s-wave scattering unitarity limit \( \Delta_{\rho_{\text{res}}}=(x/c)^{2} \) for \( n=0.15-\text{const} \), depicted as thin solid line in Fig. 2(b). Here, \( d=c/2 \approx 6 \) Å is the average separation of the CuO second layers in the La\(_{1.85}\)Sr\(_{0.15}\)Cu\(_{1-y}\)Ni\(_{y}\)O\(_{4}\) films 31. To find the actual \( n(y) \) dependence we carried out the thermopower measurements.

Ni doping increases the positive Seebeck coefficient \( S \), as can be seen in Fig. 3. Taking into account the universal correlation between \( S(290 \text{ K}) \) and hole concentration fulfilled in most of cuprate families 36-38, this strongly suggests decreasing of carrier density with \( y \). The \( S(T) \) curves in La\(_{1.85}\)Sr\(_{0.15}\)Cu\(_{1-y}\)Ni\(_{y}\)O\(_{4}\) retain the specific features of thermopower in underdoped cuprates: the initial strong growth of \( S \) with increasing \( T \) is followed by a broad maximum and subsequent slight decrease in \( S \) 39. We find that the phenomenological asymmetrical narrow-band model 40 describes the experimental \( S(T) \) curves very well at high \( T \), above their maximum at \( T_{\text{max}} \). In this model, a sharp density-of-states peak with the effective bandwidth \( W_{D} \) is located near the Fermi level \( E_{F} \) and the carriers from the energy interval \( W_{F} \) are responsible for conduction. In addition, a shift \( bW_{D} \) between the centers of \( W_{F} \) and \( W_{F} \) bands is assumed. The best fits of the formula determining \( S \) in the model (Eq. (1) in Ref. 40) are shown as thick lines in Fig. 3. The discrepancies at low \( T \) come from the limitations of the model derived under the assumption \( W_{D} \approx k_{B}T_{\text{max}} \). Above \( y=0.15 \), the model also fails for larger \( T \), well above \( T_{\text{max}} \) (\( W_{D} > 380 \text{ meV} \), while \( k_{B}T \approx 26 \text{ meV} \) at 300 K). For \( y=0.17 \) and \( y=0.19 \), \( S \) changes as \( x_{1}/T_{\text{ave}} \approx 250 \text{ K} \), consistent with the formula \( S(T)=k_{B}E_{0}/(k_{B}T+\text{const}) \) indicative of polarons. The thermal activation energy \( E_{\text{a}} \) estimated from the best fit for \( y=0.19 \), \( E_{\text{a}}=32.4 \pm 0.2 \text{ meV} \), is in good agreement with Ref. 41.

At region of interest corresponding to high \( T \), the asymmetrical narrow-band model works very well. The obtained asymmetry factor \( b = -(0.06-0.08) \), although very small, is essential for good fits. The ratio \( F \) of \( n \) to the total number od states \( n_{DOS} \) is slightly above half-filling (0.51-0.53) and thus consistent with the sign of \( S \) and in good agreement with literature data for \( x=0.15 \) LSCO 40-42. The Fermi level, \( E_{F}=(F-0.5)W_{D}-bW_{D} \),
crosses the conduction-band upper edge at \( y \approx 0.15 \). The band-filling \( F \) does not show any obvious \( y \)-dependence. This means that Ni does not change \( n \) in the system (provided \( n_{\text{DOS}} \) remains constant). The primary effect of Ni doping appears to be localization of existing mobile carriers, as revealed by decreasing of \( W_a/W_D \) ratio with increasing \( y \) (inset to Fig. 3). The ratio extrapolates to zero at \( y=0.22 \pm 0.02 \), resulting in the average “localization rate” \( \Delta n/\Delta y=0.7 \pm 0.1 \) hole/Ni ion. Employing the simple two-band model with \( T \)-linear term \( \tilde{\Gamma} \), where half-width of the resonance peak \( \Gamma \) corresponds to the range of delocalized states, results in the similar physical picture. We found that \( \Gamma \) starts to decrease with increasing \( y \) above 0.07 and approaches zero for \( y=0.17 \) (Ref. [3]). The results are consistent with measurements of local distortion around Ni ions suggesting trapping hole by each \( \text{Ni}^{2+} \) [44] to create a well-localized Zhang-Rice doublet state [45], albeit indicate a slightly lower \( \Delta n/\Delta y \).

Having established the effective mobile carrier concentration \( n \approx 0.15 - y \), we can revert to Fig. 2b. As indicated by thick solid line, scattering in the samples with the smallest \( \rho_{\text{res}} \) is in the unitarity limit for nonmagnetic impurity. Even assuming \( \Delta n/\Delta y=0.7 \) (and unitarity limit), at most only 20\% of increase in \( \rho_{\text{res}} \) for \( y=0.08 \) can be attributed to scattering on magnetic moments. Thus, while our previous finding of spin-glass behavior in the system [30] undoubtedly indicates that the magnetic role of Ni ions in the spin-1/2 network of the CuO\(_2\) planes cannot be neglected, the scattering on Ni has a predominantly nonmagnetic origin [28].

In Fig. 2c we show that the fitted \( \rho_{\text{sat}} \) agrees unexpectedly with \( \rho_{\text{IRM}} \) calculated from Boltzmann theory for the small Fermi surface with \( n \) holes. Assuming a cylindrical surface with the height \( 2\pi/d \) and taking Ioffe-Regel-Mott condition as \( l_{\text{min}} \approx a \) (i.e. \( k_F l_{\text{min}} \approx \pi \), see Ref. [46, 47]), where \( a \) is the lattice parameter in CuO\(_2\) plane, one gets \( \rho_{\text{IRM}}^{\text{small}}=\left(\sqrt{2\pi\hbar/e^2}\right)/\sqrt{n}=0.08/\sqrt{n} \) m\( \Omega \) cm [29]. This is clearly distinguishable from the large Fermi surface case, \( \rho_{\text{IRM}}^{\text{large}}=0.68/\sqrt{1+n} \), inapplicable to the system. A simple formal statistics for all 32 measured samples shows that the product \( \rho_{\text{sat}}/\sqrt{n} \) for \( n=0.15-0.02 \) has a distribution with mean \( \approx \) median and zero skewness [48].

The precise location of the large-to-small Fermi surface transition on the phase diagram of cuprates is still under debate. In the bismuth-based family, the expected linear relationship between \( n \) estimated from \( T_C \) and from Luttinger count is obtained only assuming large surface from overdoped specimens down to \( p=0.145 \) inclusive [40]. In La\(_{1.85}\)Sr\(_{0.15}\)Cu\(_{1-y}\)Ni\(_{y}\)O\(_4\), Ni doping effectively moves the system towards smaller \( p \) but the smooth evolution of all the fitted \( \rho(T) \) parameters down to \( y=0 \) points toward small Fermi surface at \( p=0.15 \). This finding is consistent with recent Hall measurements in YBa\(_2\)Cu\(_3\)O\(_y\) indicating that Fermi-surface reconstruction with decreasing doping ends sharply at \( p=0.16 \) (Ref. [16]).

The \( T \)-linear coefficient in Eq. (11) for three \( y=0 \) samples \( \alpha_1=0.93 \pm 1.0 \) \( \mu\Omega/cm/K \) is identical as that at \( n_{\text{cr}}=0.185 \pm 0.005 \) where a change in LSCO transport coefficients was found when tracked from the overdoped side [35]. Evidently, \( \alpha_1 \) is not sensitive to disappearing of the antidonal regions during degradation/reconstruction of large Fermi surface into arcs/pockets. The linear-in-\( T \) scattering is anisotropic in CuO\(_2\) plane [50] and its maximal value at \( (\pi,0) \) for \( \alpha_1=1 \) \( \mu\Omega/cm/K \) is comparable [51] with Planckian dissipation limit [51, 52]. De-coherence of quasiparticle states beyond this limit seems to be plausible explanation [53] of negligible role of antidonal states in the conductivity. While \( \alpha_1 \) vanishes for all specimens with \( y \geq 0.035 \), the \( T^2 \)-coefficient \( \alpha_2 \) changes linearly with \( 1/n \) (compare Ref. [3]) in the whole studied \( y \) range (Fig. 2b). When extrapolated outside accessible \( n \), \( \alpha_2 \) approaches zero at \( n_{\text{cr}}=0.167 \pm 0.009 \). With such estimated error, the result means that strictly linear \( \rho(T) \) dependence (albeit masked by \( \rho_{\text{sat}} \)) in LSCO is observed from the underdoped side only at optimum doping. Interpreting this as an indication of (antiferromagnetic) Quantum Critical Point, remains speculative since \( \alpha_2 \) diverges at such a point [53, 54].

After using the parallel-resistor formalism in the zero-field transport description, we will employ it to analysis of magnetoresistance in the sections below. In the strongly overdoped cuprates magnetoresistance obeys Kohler’s rule [6]. The relative change of resistivity in magnetic field \( B \), \( \Delta \rho/\rho_0 \), is a unique function of \( B/\rho_0 \), where \( \rho_0 \equiv \rho(B=0) \). In recent re-analysis of \( x=0.09 \) LSCO data from Ref. [9], the modified Kohler’s rule was proposed [4]. The isotherm temperature transverse magnetoresistance vs. \( B/\rho_0 \) curves appeared to collapse onto single curve when \( \rho_0 \) is replaced by \( \rho_{\text{sat}}/\rho_{\text{res}} \). Here, we propose an alternative approach for considering the large LSCO resistivity in the magnetoresistance analysis.
At the lowest temperatures down to 2 K, all non-superconducting samples exhibit large and negative in-plane (T || ab) magnetoresistance, both in the transverse (TMR, B || ab) and longitudinal (LMR, B || ab) configurations [TMR(B=9T)≈5×10⁻² at 2 K]. In the following, we focus on the high-T region where, for the whole y range studied, magnetoresistance in both configurations is positive and two orders of magnitude smaller than in the low-T region. The typical field dependence of resistivity at various temperatures is displayed for y=0.06 specimen in the right inset to Fig. 4(a). Above 35 K, ρ increases as B². Below 60 K, the positive TMR begins to decrease with decreasing T and smoothly evolves into negative one at low T. Similar behavior is observed for LMR, which constitutes a significant portion of TMR (being equal to 60% of TMR at T=150 K as an example) and thus may not be ignored in the analysis. To obtain the orbital part, OMR, at T \( \geq 35 \) K, we fitted ρ(B) with the form \[ \Delta \rho(B)/\rho(0)\] TMR,LMR = \( a_{\text{TMR,LMR}} B^2 \) and next calculated \( a_{\text{ORB}} = \Delta \rho_{\text{ORB}}/\rho_{\text{ORB}} \). The extracted coefficients are displayed in Fig. 4(a). A reliable and precise comparison of the various possible OMR scaling requires moderate numerical smoothing without alternating the \( a_{\text{ORB}} \) vs. T dependence [left inset to Fig. 4(a)]. Employing the smoothed \( a_{\text{ORB}} \) coefficients, \( a_{\text{orb}} \), OMR at any field B can be calculated as \( \text{OMR}=a_{\text{orb}} B^2 \).

Clearly, Kohler’s rule in La\(_{1.85}\)Sr\(_{0.15}\)Cu\(_{1-y}\)Ni\(_y\)O\(_4\) is violated when applied directly to the measured OMR of the specimen [inset to Fig. 4(c)]. Modification of the rule in the way described in Ref. 4 does not lead to any reasonable scaling range. The \( \Delta \rho_{\text{orb}}/(\rho_0 - \rho_{\text{res}}) \) vs. \( B/(\rho_0 - \rho_{\text{res}}) \) lines collapse one onto another between 180 K and \( T_{\text{up}} = 200 \) K, spanning only 10% of \( T_{\text{up}} \). Let’s note that the existence of such a scaling - where the whole absolute resistivity change in field, \( \Delta \rho_{\text{orb}} \), is related only to T-dependent \( \rho_{\text{sat}}-\rho_{\text{res}} \) part - would mean in the classical picture that the field acts between the scattering events only on these carriers that scatter against phonons during the subsequent scattering event and does not bend trajectories of those scattered against impurities.

The OMR analysis reveals that Kohler’s rule in non-superconducting specimens is fulfilled in the ideal branch of the parallel-resistor model where the influence of the shunt \( \rho_{\text{sh}} \) is eliminated [Fig. 4(b)]. Interpreting the model in the spirit of the minimal \( \tau_0 \) leads to the assumption that \( \rho_{\text{sat}} \) is field-independent, at least in the weak-field regime (where actually observed OMR \( \times B^2 \) dependence is expected). Fitting of Eq. 1 \( \rho(T) \) measured at various fields up to 9T does not reveal any systematic change of \( \rho_{\text{sat}} \) with \( B \). With \( \rho_{\text{sat}}(B)≈\rho_{\text{sat}}(0) \), OMR of the ideal branch, \( \Delta \rho_{\text{id}}/\rho_{\text{id}}(0) \) can be calculated from the measured quantities employing only one fitting parameter \( \rho_{\text{sat}}(0) \). The obtained \( \Delta \rho_{\text{id}}/\rho_{\text{id}}(0) \) vs. \( B/(\rho_{\text{id}}(0)) \) curves from 110 K up to \( T_{\text{up}} = 210 \) K collapse to a single temperature-independent line, spanning 50% of \( T_{\text{up}} \). The similar result was obtained for y=0.07 specimen. Closer to the superconducting region of the phase diagram, for y=0.04, the scaling interval in \( \rho_{\text{sat}}(B,T) \) is much smaller (140K-160K) but larger difference between \( \rho_{\text{sat}} \) and \( \rho_{\text{res}} \) emphasizes the difference between the possible OMR scalings and the scaling approach illustrated in Fig. 4(c) fails completely.

Concluding, the signs of resistivity saturation both at the value \( \rho_{\text{IRM}} \) and with n-dependence from Boltzmann theory reflect metallic-like character of transport despite small volume of Fermi surface and strong disorder in underdoped LSCO. A fully quantum-mechanical explanation of saturation is still lacking. Evidently however, strong electron-electron interactions do not invalidate the \( \rho_{\text{IRM}} \) limit. The Ni-induced order-of-magnitude increase of \( \rho_0/\rho_{\text{sat}} \) ratio leaves the resulting \( \rho_{\text{IRM}} \) intact. The revealed omnipresence of \( \tau_{\text{min}} \) even in bad metals points towards universality of the Ioffe-Regel-Mott criterion rather than accidental-only agreement between \( \rho_{\text{IRM}} \) and saturation.

The known huge increase of \( \rho \) in LSCO outside our \( T \)-measurement window means the breakdown of quasiparticle description. At lower \( T \), the transport in...
La$_{1.85}$Sr$_{0.15}$Cu$_{1-y}$Ni$_y$O$_4$ has a coherent description in the framework of Fermi-liquid theory, where Kohler’s rule is derived under single-relaxation-time approximation with the assumption of small $\tau$-anisotropy over Fermi surface [63, 65]. Fulfillment of the rule when $\rho_0$ is changed by changing temperature indicates nearly $T$-independent frequency distribution of the phonons involved [65] and is consistent with Fermi-arc length in cuprates being constant [64] rather than decreasing with decreasing temperature [66, 67]. Ni doping can restore antiferromagnetic fluctuations [68] that give additional $T$-dependence in nearly-antiferromagnetic-metals magnetoresistance via correlation length $\xi_{AF}(T)$, OMR $\propto \xi_{AF}^2(T)/\rho_0^2$ (Ref. [63]). Thus fulfillment of conventional Kohler’s rule suggests $T$-independent $\xi_{AF}$ within the framework of this theory.

In summary, the typical for normal metals parallel-resistor formalism, employed for analysis of La$_{1.85}$Sr$_{0.15}$Cu$_{1-y}$Ni$_y$O$_4$ transport, reveals resistivity saturation at the value expected from Boltzmann theory for the small Fermi surface and fulfillment of unmodified Kohler’s rule in the nonsuperconducting part of the phase diagram.

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[1] S. E. Sebastian, N. Harrison, M. M. Altarawneh, R. Liang, D. A. Bonn, W. N. Hardy, and G. G. Lonzarich, Phys. Rev. B 81, 140505 (2010)
[2] S. I. Mirzaei, D. Stricker, J. N. Hancock, C. Berthod, A. Georges, E. van Humen, M. K. Chan, X. Zhao, Y. Li, M. Greven, et al., Proceedings of the National Academy of Sciences 110, 5774 (2013)
[3] N. Barišić, M. K. Chan, Y. Li, G. Yu, X. Zhao, M. Dressel, A. Smontara, and M. Greven, Proceedings of the National Academy of Sciences 110, 12235 (2013)
[4] M. K. Chan, M. J. Veit, C. J. Dorow, Y. Ge, Y. Li, W. Tabis, Y. Tang, X. Zhao, N. Barišić, and M. Greven, Phys. Rev. Lett. 113, 177005 (2014)
[5] H. Esiński, N. Kaneko, D. L. Feng, A. Damascelli, P. K. Mang, K. M. Shen, Z.-X. Shen, and M. Greven, Phys. Rev. B 69, 064512 (2004)
[6] Y. Ando, G. S. Boebinger, A. Passner, N. L. Wang, C. Geibel, and F. Steglich, Phys. Rev. Lett. 77, 2065 (1996)
[7] Y. Ando, Y. Kurita, S. Komiya, S. Ono, and K. Segawa, Phys. Rev. B 92, 197001 (2014)
[8] J. M. Harris, Y. F. Yan, P. Matl, N. P. Ong, P. W. Anderson, T. Kimura, and K. Kitazawa, Phys. Rev. Lett. 75, 1391 (1995)
[9] T. Kimura, S. Miyasaka, H. Takagi, K. Tamasaku, H. Esiński, S. Uchida, K. Kitazawa, M. Hiroi, M. Sera, and N. Kobayashi, Phys. Rev. B 53, 8733 (1996)
[10] R. H. McKenzie, J. S. Qualls, S. Y. Han, and J. S. Brooks, Phys. Rev. B 57, 11854 (1998)
[11] M. Platé, J. D. F. Mottershead, I. S. Elffimov, D. C. Peets, R. Liang, D. A. Bonn, W. N. Hardy, S. Chiuzbaian, M. Falub, M. Shi, et al., Phys. Rev. Lett. 95, 077001 (2005)
[12] B. Vignolle, A. Carrington, R. A. Cooper, M. M. J. French, A. P. Mackenzie, C. Jaouet, D. Vignolles, C. Proust, and N. E. Hussey, Nature 455, 952 (2008)
[13] B. Vignolle, D. Vignolles, D. LeBoeuf, S. Lepault, B. Ramshaw, R. Liang, D. Bonn, W. Hardy, N. Doiron-Leyraud, A. Carrington, et al., C. R. Phys. 12, 446 (2011)
[14] N. Doiron-Leyraud, C. Proust, D. LeBoeuf, J. Levallois, J.-B. Bonnemaison, R. Liang, D. A. Bonn, W. N. Hardy, and L. Taillefer, Nature 447, 565 (2007)
[15] N. Barišić, S. Badoux, M. K. Chan, C. Dorow, W. Tabis, B. Vignolle, G. Yu, J. Beard, X. Zhao, C. Proust, et al., Nat. Phys. 9, 761 (2013)
[16] S. Badoux, W. Tabis, F. Laliberté, G. Grissonnanche, B. Vignolle, D. Vignolles, J. Béard, D. A. Bonn, W. N. Hardy, R. Liang, et al., Nature 531, 210 (2016)
[17] S. Badoux, S. A. A. Afshar, B. Michon, A. Ouellet, S. Fortier, D. LeBoeuf, T. P. Croft, C. Lester, S. M. Hayden, H. Takagi, et al., Phys. Rev. X 6, 021004 (2016)
[18] N. F. Mott, Metal-Insulator Transition (Taylor & Francis Ltd., London, 1990), chap. 1, p. 30, 2nd ed.
[19] H. Wiesmann, M. Gurvitch, H. Lutz, A. Ghosh, B. Schwarz, M. Strongin, P. B. Allen, and J. W. Halley, Phys. Rev. Lett. 38, 782 (1977)
[20] J. Orenstein, G. A. Thomas, A. J. Millis, S. L. Cooper, D. H. Rapkine, T. Timusk, L. F. Schneemeyer, and J. V. Waszczak, Phys. Rev. B 42, 6342 (1990)
[21] H. Takagi, B. Batlogg, H. L. Kao, J. Kwo, R. J. Cava, J. J. Krajewski, and W. F. Peck, Phys. Rev. Lett. 69, 2975 (1992)
[22] N. L. Wang, B. Buschinger, C. Geibel, and F. Steglich, Phys. Rev. B 54, 7445 (1996)
[23] K. Takenaka, J. Nohara, R. Shiozaki, and S. Sugai, Phys. Rev. B 68, 134501 (2003)
[24] H. Takagi, T. Ido, S. Ishibashi, M. Ueta, S. Uchida, and Y. Tokura, Phys. Rev. B 40, 2254 (1989)
[25] K. Takenaka, R. Shiozaki, S. Okuyama, J. Nohara, A. Osuka, Y. Takayanagi, and S. Sugai, Phys. Rev. B 65, 092405 (2002)
[26] O. Gunnarsson, M. Calandra, and J. E. Han, Rev. Mod. Phys. 75, 1085 (2003)
[27] M. Calandra and O. Gunnarsson, Europhys. Lett. 61, 88 (2003)
[28] E. W. Hudson, K. M. Lang, V. Madhavan, S. H. Pan, H. Esiński, S. Uchida, and J. C. Davis, Nature 411, 920 (2001)
[29] Y. Tanabe, T. Adachi, K. Suzuki, T. Kawamata, Risdiyan, A. Malinowski, V. L. Bezusyy, R. Minikayev, P. Dziawa, Y. Syryanyy, and M. Sawicki, Phys. Rev. B 84, 024409 (2011)
[30] A. Malinowski et al. (unpublished)
[32] A. Carrington and J. Cooper, Physica C: Superconductivity 219, 119 (1994)
[33] M. Gurvitch, Phys. Rev. B 24, 7404 (1981)
[34] Hussey, N. E., Eur. Phys. J. B 31, 495 (2003)
[35] R. A. Cooper, Y. Wang, B. Vignolle, O. J. Lipscombe, S. M. Hayden, Y. Tanabe, T. Adachi, Y. Koike, M. No- 
hara, H. Takagi, et al., Science (2009)
[36] S. D. Obertelli, J. R. Cooper, and J. L. Tallon, Phys. Rev. B 46, 14928 (1992)
[37] T. Honma, P. H. Hor, H. H. Hsieh, and M. Tanimoto, Phys. Rev. B 70, 214517 (2004)
[38] M. Gurvitch, Phys. Rev. B 32, 15116 (1985)
[39] J. S. Wen, Z. J. Xu, G. Gu, and A. Kaminski, Phys. Rev. B 37, 3759 (1988)
[40] N. E. Hussey, K. Takenaka, and H. Takagi, Phil. Mag. 84, 2847 (2004)
[41] M. R. Graham, C. J. Adkings, H. Behar, and R. Rosen- baum, J. Phys.: Condensed Matter 10, 809 (1998)
[42] The linear fit to $\rho_{sat} \sqrt{m}(y)$ for $n = 0.15 - y$ has a small nonzero slope that vanishes for $n = 0.15 - 0.7y$, but the result should be taken with some caution because of rel- atively large scattering of data.
[43] Y. He, Y. Yin, M. Zech, A. Soumyanarayanan, M. M. Yee, T. Williams, M. C. Boyer, K. Chatterjee, W. D. Wise, I. Zeljkovic, et al., Science 344, 608 (2014)
[44] M. Abdel-Jawad, M. P. Kennett, L. Balicas, A. Carring- 
ton, A. P. Mackenzie, R. H. McKenzie, and N. E. Hussey, Nat Phys 2, 821 (2006)
[45] J. Zaanen, Nature 430, 512 (2004)
[46] C. C. Homes, S. V. Dordevic, M. Strongin, D. A. Bonn, R. Liang, W. N. Hardy, S. Komiyi, Y. Ando, G. Yu, N. Kaneko, et al., Nature 430, 539 (2004)
[47] S. A. Grigera, R. S. Perry, A. J. Schofield, M. Chiao, S. R. Julian, G. G. Lonzarich, S. I. Ikeda, Y. Maeno, A. J. Millis, and A. P. Mackenzie, Science 294, 329 (2001)
[48] P. Gegenwart, J. Custers, C. Geibel, K. Neumaier, T. Tayama, K. Tenya, O. Trovarelli, and F. Steglich, Phys. Rev. Lett. 89, 056402 (2002)
[49] J. Paglione, M. A. Tanatar, D. G. Hawthorn, E. Boaknin, R. W. Hill, F. Ronning, M. Sutherland, and L. Taillefer, Phys. Rev. Lett. 91, 246405 (2003)
[50] A. Malinowski, M. F. Hundley, C. Capan, F. Ronning, R. Movshovich, N. O. Moreno, J. L. Sarrao, and J. D. Thompson, Phys. Rev. B 72, 184506 (2005)
[51] The $T$-range of the scaling has been found semiquantita- tively as a maximal one for which adding a next measured $T$-point outside this range doesn’t broaden the existing angle between the outermost lines in the plot.
[52] The larger field is needed to unmask the possible $B$-dependence of $\rho_{sat}$ (when treated as a formal fitting pa- rameter) that may be obscured inside small magnetore- sistance of the system.
[53] M. Calandra and O. Gunnarsson, Phys. Rev. B 66, 205105 (2002)
[54] P. B. Allen, Physica B: Condensed Matter 318, 24 (2002)
[55] Y. Werman and E. Berg, Phys. Rev. B 93, 075109 (2016)
[56] Even for large $y$, the electron-electron scattering $aT^2$ term at $T=350$ K exceeds the $T$-dependent Ni-induced $a_0$ term.
[57] J. M. Ziman, Principles of the Theory of Solids (Cambridge University Press, Cambridge, 1972), chap. 7, 2nd ed.
[58] H. Kontani, Reports on Progress in Physics 71, 026501 (2008)
[59] A. B. Pippard, Magnetoresistance in Metals (Cambridge University Press, Cambridge, 1989), p. 23
[60] T. Kondo, A. D. Palczewski, Y. Hamaya, T. Takeuchi, J. S. Wen, Z. J. Xu, G. Gu, and A. Kaminski, Phys. Rev. Lett. 111, 157003 (2013)
[61] A. Kaminski, T. Kondo, T. Takeuchi, and G. Gu, Philo- sophical Magazine 95, 453 (2015)
[62] A. Kanigel, M. R. Norman, M. Randera, U. Chatterjee, S. Souma, A. Kaminski, H. M. Fretwell, S. Rosenkranz, M. Shi, T. Sato, et al., Nat Phys 2, 447 (2006), ISSN 1745-2473
[63] K. Tsutsui, A. Toyama, T. Tohyama, and S. Maekawa, Phys. Rev. B 80, 224519 (2009)