Metal-Insulator Crossover in Superconducting Cuprates in Strong Magnetic Fields

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The metal-insulator crossover of the in-plane resistivity upon temperature decrease, recently observed in several classes of cuprate superconductors, when a strong magnetic field suppresses the superconductivity, is explained using the U(1) $\times$ SU(2) Chern-Simons gauge field theory. The origin of this crossover is the same as that for a similar phenomenon observed in heavily underdoped cuprates without magnetic field. It is due to the interplay between the diffusive motion of the charge carriers and the "peculiar" localization effect due to short-range antiferromagnetic order. We also calculate the in-plane transverse magnetoresistance which is in fairly good agreement with available experimental data.

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The in-plane resistivity in heavily underdoped samples of cuprates [in particular, $La_{2-x}Sr_xCuO_4$ (LSCO)] exhibits a minimum and a crossover from metallic to insulating behavior upon the temperature decrease [1]. Recently a similar crossover was observed in several classes of superconducting cuprates [2–5] when a strong magnetic field (up to 60 T) suppresses the superconductivity. The "obvious" interpretation in terms of two-dimensional (2D) localization or 2D insulator-superconductor transition is ruled out, as the sheet resistance, defined as $\rho_{sh} \equiv \rho_{ab}/a$, where a is the interlayer distance, at the crossover point is between 1/25 to 1/12 in units of h/e^2 [2–4], or, using a free electron model, the estimated product $k_F l$, where k_F is the Fermi momentum and l is the mean free path, is between 12 and 25. The insulating ground state persists up to optimal doping in LSCO [2] and in electron-doped superconductors $Pr_{2-x}Ce_xCuO_4$ [3], while in newly studied $Bi_2Sr_{2-x}La_xCuO_{6+\delta}$ (La-doped Bi-2201) it persists only up to 1/8 hole-doping without showing any signature of stripe formation. Thus ascribing this metal-insulator (MI) crossover to a quantum critical point related to charge density instability [6] is open to objections [7]. There were several attempts to interpret the insulating ground state in doped cuprates using various arguments on non-Fermi liquid behavior [8]. However, the crossover phenomenon as temperature varies has not been addressed up to now, to the best of our knowledge. It was thought earlier that the MI crossover in the absence of magnetic field [1] and that induced by magnetic field is of different origin (the sheet resistance in the first case was substantially higher) [2]. This is doubtful, because the recent experiments on $YBa_2Cu_3O_{7-y}$ (YBCO) show a MI crossover in nonsuperconducting compounds in the absence of magnetic field [9], and the same type of MI crossover (with comparable sheet resistance at the crossover) in the superconducting compounds of the same series doped with Zn in the presence of magnetic field [5]. We believe the *two crossover phenomena are of the same origin*. We have used the U(1) $\times$ SU(2) Chern-Simons (CS) approach to the t-J model, proposed by us earlier [10] to explain the MI crossover in heavily underdoped cuprates in the absence of magnetic field [11]. In this Letter we generalize our formalism to include the effect of magnetic field and show that such a MI crossover is a universal feature of doped cuprates, and it is due to a "peculiar charge localization" effect (using the wording of Ref. [4]), resulting from the interplay of the spin-excitation gap [corresponding to short-range antiferromagnetic order (SRAFO)] and the holon induced anomalous dissipation. Moreover, we show that the observed large positive in-plane transverse magnetoresistance (MR) at low temperatures [12,13] can be semiquantitatively explained within this formalism.

The U(1) $\times$ SU(2) CS gauge field approach is a particular scheme of slave-particle formalism to treat the t-J model based on an exact identity [14], introducing a U(1) field gauging the global charge symmetry and a SU(2) field gauging the global spin symmetry, both with CS actions. Using an optimization procedure of free energy [10], a careful mean field (MF) approximation gives the following results: The U(1) gauge field for low doping δ develops a π flux per plaquette converting holons into Dirac fermions with a Fermi energy $\epsilon_F \sim t\delta$. The holons induce a vortex structure in the MF configurations of the SU(2) gauge field, with vortices centered at the holon positions. These dressed holons in turn are seen as slowly moving impurities by spin waves giving rise to a spinon mass $m_s \sim \sqrt{\delta|\ln\delta|}$. Notice that this feedback is self-consistent because for low δ we have $\epsilon_F \sim t\delta \ll \epsilon_s \sim J\sqrt{\delta|\ln\delta|}$. We use $J = 0.1$ eV, $t/J = 3$ in our numerical computations. Because of a special choice of "gauge fixing" (using the Néel gauge) the SU(2) gauge field becomes physical, describing the spin fluctuations. The spinon action is given by a nonlinear σ model with a mass term (spinon gap) which in the CP^1 representation yields a new self-generated U(1) gauge field A coupling holons and spinons [this field is analogous to the U(1) gauge field in the standard slave-particle approaches [15,16]]. Because of coupling to holons (fermions in our approach), this gauge field acquires an anomalous dissipation term, the "Reizer singularity" [17], which dominates the low-energy action for the transverse component of the gauge field A^T. For $\omega, |\vec{q}|, \omega/|\vec{q}| \sim 0$ we have $\langle A^T A^T \rangle(\omega, \vec{q}) \sim (\chi |\vec{q}|^2 + i\kappa \frac{\omega}{|\vec{q}|})^{-1}$, where $\chi \sim t/\delta$ is the diamagnetic susceptibility and $\kappa \sim \delta$ is the Landau damping. The interplay of the two different energy scales, the spinon gap and the holon induced anomalous dissipation, is the key factor in our interpretation of the MI crossover [11]. For the temporal component in the same limit, we have $\langle A^0 A^0 \rangle(\omega, \vec{q}) \sim (\gamma + \omega_p)^{-1}$, where γ is the fermion density of states and ω_p is the plasmon gap.

Now we consider the introduction of a magnetic field H perpendicular to the plane. The Ioffe-Larkin rule [18], $R = R_s + R_h$, i.e., the observed resistivity is the sum of the holon and spinon contributions, can be generalized to this case. The external electromagnetic potential $A_{\mathrm{e.m.}}$, corresponding to the constant magnetic field H, can couple with coefficient $-\varepsilon$ to spinons and $1 - \varepsilon$ to holons. In principle, $0 \leq \varepsilon \leq 1$ is arbitrary. However, to be consistent with the requirement that the physical inverse magnetic susceptibility should be the sum of the inverse of that of holons and spinons, i.e., $\chi^{-1} = (\chi_s^*)^{-1} + (\chi_h^*)^{-1}$, where χ_s^* and χ_h^* are the renormalized spinon and holon susceptibility, respectively (this relation can be derived in the same way as the Ioffe-Larkin rule), we find $\varepsilon = \chi_h^*/(\chi_h^* + \chi_s^*)$. This value was argued earlier using variational considerations [19]. Replacing these quantities by unrenormalized values, we find $1 - \varepsilon \sim \chi_s/(\chi_h + \chi_s) \sim \frac{J}{t}\sqrt{\delta/|\ln \delta|} \ll 1$. In the Coulomb gauge $A^0_{\mathrm{e.m.}} = 0$, the effective action for the gauge field A can be written as

$$S_{\mathrm{eff}}(A) = \int dx^0\, d^2x \left[\frac{i}{2} \{ A^0(\Pi_h^0 + \Pi_s^0)A^0 + [A^T + (1-\varepsilon)A_{\mathrm{e.m.}}]\Pi_h^{\perp}[A^T + (1-\varepsilon)A_{\mathrm{e.m.}}] + (A^T - \varepsilon A_{\mathrm{e.m.}})\Pi_s^{\perp}(A^T - \varepsilon A_{\mathrm{e.m.}}) \} + \frac{i\sigma_h(H)}{2\pi} A^0 \epsilon_{ij} \partial^i A^j \right], \tag{1}$$

where $\Pi_h^{0,\perp}, \Pi_s^{0,\perp}$ are unrenormalized polarization bubbles due to holons and spinons, respectively, and $\sigma_h(H)$ is the Hall conductivity due to holons. Note $A_{\mathrm{e.m.}}$ appears in two places in this low-energy effective action: one is simply a shift of the transverse component of the gauge field A^T by $(1-\varepsilon)A_{\mathrm{e.m.}}$ and $-\varepsilon A_{\mathrm{e.m.}}$ corresponding to the minimal coupling to holons and spinons, respectively, while the other is a CS term due to parity breaking induced by H.

As remarked in [19], the leading effect of the integration over A_0 is the renormalization of the diamagnetic susceptibility: $\chi \to \chi(H) = \chi + \frac{\sigma_h^2(H)}{4\pi^2\gamma}$ in the A^T effective action. The holon contribution R_h can be evaluated using the Boltzmann equation, taking into account the classical cyclotron effect, as in [19], obtaining

$$R_h = R_h^0 \left[1 + \left(\frac{(1-\varepsilon)H\tau}{m_h} \right)^2 \right], \qquad R_h^0 \sim \frac{m_h}{8}\frac{1}{\tau} \sim \delta \left[\frac{1}{\epsilon_F \tau_{\mathrm{imp}}} + \left(\frac{T}{\epsilon_F} \right)^{4/3} \right], \tag{2}$$

where τ is the transport relaxation time, τ_{imp} is the impurity scattering time, and $m_h \sim \delta/t$ is the holon mass. The spinon contribution R_s is evaluated here using the Kubo formula for the spinon current: $R_s = \lim_{\omega \to 0} \omega [\mathrm{Im}\Pi_s^{\perp}(\omega)]^{-1}$, where $\Pi_s^{\perp}$ denotes the transverse polarization bubble at $\vec{q} = 0$, renormalized by gauge fluctuations. At large scales, for $x^0 \gg |\vec{x}|$, $\Pi_s^{\perp}(x)$ is approximately given by $\langle \partial_\mu G(x|A) \partial^\mu G(x|A) \rangle$, where $\langle\,\rangle$ denotes the A-expectation value and $G(x|A)$ is the spinon propagator. Using the Fradkin representation [11,20] it can be transformed into a gauge invariant form $\langle \partial_\mu G(x|F) \partial^\mu G(x|F) \rangle$, where $G(x|F)$ in terms of a path integral over 3-velocities, $\phi^\mu(t) \equiv \dot{q}_\mu(t)$, $\mu = 0, 1, 2$, is given by

$$G(x|F) \sim i \int_0^\infty ds\, e^{-im_s^2 s} \int \mathcal{D}\phi^\mu\, e^{\frac{i}{4}\int_0^s \phi_\mu^2(t)\, dt} \times \int d^3p\, e^{ipx - ip^2 s} \times e^{i\mathcal{Q}^{ij}(p,s,\phi|s',\lambda)[F_{ij}(p,s|s',\lambda) + \epsilon_{ij}\varepsilon H]} \tag{3}$$

with

$$\mathcal{Q}^{ij}(p,s,\phi|s',\lambda)[\#] = \int_0^1 d\lambda\, \lambda \int_0^s ds' \int_0^{s'} ds'' \times [\phi^i(s') - 2p^i] \times [\phi^j(s'') - 2p^j][\#] \tag{4}$$

and

$$F_{ij}(p,s|s',\lambda) = F_{ij}\left(x + \lambda \int_0^{s'} ds''\, \phi(s'') - 2ps \right), \qquad F_{ij} \equiv \partial_i A_j - \partial_j A_i . \tag{5}$$

After integration over A^T, using the effective action, and integration over the 3-velocities in the Gaussian approximation, and over p and s by saddle point for low T, one obtains at large scales

$$\Pi_s^{\perp}(x) \sim \left[\frac{\partial}{\partial x^\mu} \left(e^{-i\sqrt{m^2 - \frac{T}{\chi(H)} f(\alpha) + \frac{\alpha^2 \varepsilon^2 H^2}{3q_0^2}}\sqrt{x_0^2 - |\vec{x}|^2}} \times e^{-\frac{T}{4\chi(H)} q_0^2 g(\alpha) \frac{(x_0^2 - |\vec{x}|^2)}{m^2}} \frac{1}{\sqrt{(x_0^2 - |\vec{x}|^2)}} \right) \right]^2, \tag{6}$$

where $\alpha = \frac{q_0 |\vec{x}|}{2}$; q_0 is the momentum scale associated with the anomalous skin effect due to Reizer singularity:

$q_0 \sim (\frac{\delta^2 T}{t})^{1/3}$. f and g are functions describing the effect of gauge fluctuations and for a real argument, f is monotonically increasing, vanishing quadratically near the origin and g is monotonically decreasing vanishing at large argument. (See [21] for explicit expressions.) The integration over $|\vec{x}|$ and x^0 appearing in the Fourier transformation are evaluated by saddle point for $|\vec{x}|$, at large x^0, and with scale renormalization by principal part evaluation for x^0 [11,21]. The integrals are dominated by a complex saddle point at $|\vec{x}| = 2q_0^{-1} e^{i\frac{\pi}{4}}$ for $\chi(H) q_0 |m_s(T,H)| \lesssim T$, $\mathrm{Im}[m_s^2(T,H)] \lesssim m_s^2$, where

$$m_s^2(T,H) = m_s^2 - i\left(\frac{cT}{\chi(H)} - \frac{\varepsilon^2 H^2}{3q_0^2}\right), \tag{7}$$

with $c = -if(e^{i\pi/4})$, a constant with real part ~ 3 and a small imaginary part. For the range of physical parameters considered here ($H \lesssim 100$ T), these bounds gave a temperature range validity lying between a few tens and a few hundreds of degrees.

The saddle point produces the following effects: it induces "renormalization" of the spinon mass yielding a T and H dependent damping: $m_s^2 \to m_s^2(T,H)$; it gives rise to an attraction between spinon and antispinon leading to the formation of a damped spin wave; it introduces a multiplicative renormalization of the correlation functions which for R_s is given by $Z(T,H)[m_s^2(T,H)]^{\frac{1}{8}}$, where $Z(T,H) = (c' \frac{T}{\chi(H)} q_0^{-3} - \frac{2}{3}\varepsilon^2 H^2 q_0^{-5})^{1/2}$ with c' a new constant $\sim f''(e^{i\pi/4})$.

The final result in the range of T described above is given by

$$R_s \sim Z(T,H) \frac{|m_s(T,H)|^{1/4}}{\sin \frac{\Theta(T,H)}{4}}, \tag{8}$$

where $m_s(T,H) \equiv |m_s(T,H)| e^{-i\Theta(T,H)}$. The basic features of our formulas can be summarized as follows: for low T, the effect of the spinon gap is dominating ($\Theta \searrow 0$), leading to an insulating behavior; at higher temperatures one finds a metallic behavior due to the dissipation induced by gauge fluctuations, contained in $|m_s(T,H)|$, that becomes the dominant effect. Therefore a MI crossover is recovered, decreasing the temperature. The minimum of R as a function of T, $T_{MI}(\delta, H)$ is decreasing with δ and increasing with H {see the MR curve $[R(H) - R(0)]/R(0)$ in Fig. 1}, in agreement with experiment [5]. In the absence of magnetic field the crossover is determined by the interplay between $m_s^2 = \xi^{-2}$ and $T/\chi \sim Tm_h \sim \lambda^{-2}$, with ξ the magnetic correlation length and λ the thermal de Broglie wavelength. When $\lambda \le \xi$, the "peculiar" localization effect due to SRAFO is not "felt," and a metallic behavior is observed. In the opposite limit $\lambda \gg \xi$ the cuprate is insulating. The external magnetic field effectively reduces the thermal energy, or increases the thermal wavelength, so the crossover temperature goes up. The resistivity is diverging at $T = \frac{\varepsilon^2 H^2 \chi(H)}{3cq_0^2}$, which approaches $T = 0$ as H vanishes. This divergence is lying outside the region of validity of our formulas and should be considered as an artifact.

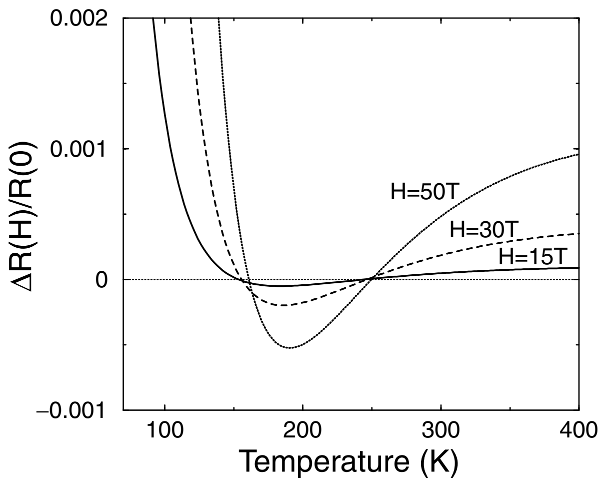

FIG. 1. The calculated magnetoresistance for cases when the quantum effects related to $\sigma_h(H)$ are strong, for doping $\delta = 0.05$. It becomes negative near the minimum which itself shifts to higher temperatures upon the field increase.

However, the shift in MI crossover temperature leads to a large positive (in-plane transverse) MR at low T which is our main new result, and it was absent in the earlier treatments [19]. The derived MR scales quadratically with H (see Fig. 2) in agreement, in particular, with data on LSCO [12,13], away from the doping $\delta = 1/8$ where the stripe effects dominate. As remarked in [19], the shift of χ induced by the CS term reduces the damping and the H^2 term due to minimal coupling acts in the same direction. In the region of T where dissipation dominates this induces a reduction of resistivity, a tendency contrasted by the classical cyclotron effect. One then has two possible types of MR curves: one is always positive but it exhibits a knee below

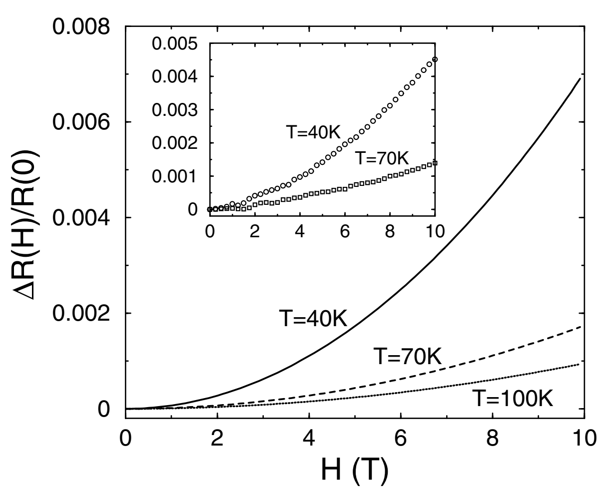

FIG. 2. The calculated field dependence of the magnetoresistance for doping $\delta = 0.075$, in comparison with experimental data on $La_{1.925}Sr_{0.075}CuO_{4+\epsilon}$ (inset), taken from Ref. [13].

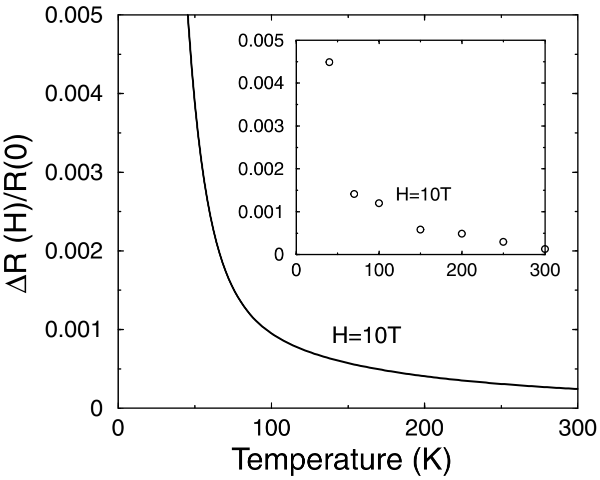

FIG. 3. The calculated temperature dependence of the magnetoresistance for doping $\delta = 0.075$, in comparison with experimental data on $La_{1.925}Sr_{0.075}CuO_{4+\epsilon}$ (inset), taken from Ref. [13].

the crossover temperature between the mass gap and the dissipation dominated regions (see Fig. 3). This behavior can be compared with the one observed in LSCO reported in [13] and we find reasonably good agreement. If, on the contrary, the quantum effects related to $\sigma_h(H)$ are sufficiently strong, a minimum develops, eventually leading to a negative MR in some region around it. This is illustrated in Fig. 1. We should point out that the large positive MR at low temperatures is foreseen in this theory for both cases.

A comment on the limit $H \sim 0$ is in order, where we recover the results of [10], in particular, $m_s^2(T,0) = m_s^2 - icT/\chi, Z(T,0) \sim \frac{1}{\sqrt{\delta}}$. In this limit the resistivity exhibits an inflection point at temperature $T^*(\delta) \sim 200\text{–}300$ K (found also in the experimental curves), above which the theoretical curves start to deviate strongly from the experimental data. We propose to interpret this inflection point as a midpoint of a crossover to a new "phase" where our MF treatment is not valid anymore. If we identify our $T^*(\delta)$ with the crossover temperature T^* found in experiments, both MI crossover temperature $T_{MI}(\delta) \equiv T_{MI}(\delta, 0)$ and $T^*(\delta)$ are found in reasonable agreement with experimental data (in the range $0.02 \lesssim \delta \lesssim 0.08$), due to a delicate cancellation of doping dependences: $\frac{T}{\chi m_s^2} \sim \frac{T\delta}{t\delta|\ln\delta|} = \frac{T}{t\ln\delta}$. R_s in this limit can be written in terms of a dimensionless variable $x \equiv cT/\chi m_s^2$ apart from an overall factor $\sqrt{|\ln\delta|}$. Hence, if we neglect the contribution $\sim T^{4/3}$ due to holons and define a "normalized resistivity" by $\tilde{R} \equiv [R - R(T_{MI})]/[R(T^*) - R(T_{MI})]$, this is a function only of x, thus exhibiting a "universal" behavior, as observed in YBCO [22].

As a final remark it might be worthwhile to notice that the same $U(1) \times SU(2)$ approach is able to reproduce qualitatively [10,23] the behavior of the spin lattice relaxation rate $(1/T_1T)^{63}$ found in underdoped YBCO [24] and a structure of Fermi arcs around $(\frac{\pi}{2}, \frac{\pi}{2})$ in the spectral density detected by ARPES [25].

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