Effect of the pseudogap on the Hall conductivity in underdoped YBa$_2$Cu$_3$O$_{6+x}$.

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In underdoped YBa$_2$Cu$_3$O$_x$ (YBCO) with $x = 6.63$, the opening of the pseudogap at $T^* \approx 160K$ has a strong effect on the Hall angle $\tan \theta$. While the Hall response is significantly reduced, the diagonal current is relatively unaffected. The Hall conductivity suppression continues deep into the flux-flow state (from $T_c$ to 40 K), as an anomalous suppression of the vortex Hall current.

72.15.Gd,74.72.Bk,74.25.Fy,74.40.+k

The gradual opening of a pseudogap affects many of the electronic properties of the underdoped cuprates, notably the nuclear magnetic resonance (NMR) relaxation rate, the electronic heat capacity, and the transport properties. Recent angle-resolved photoemission (ARPES) experiments have found that the symmetry of the pseudogap is consistent with $d$-wave symmetry. The nature of the pseudogap phase is a subject of strong current interest.

The effect of pseudogap formation on the in-plane resistivity $\rho$ was initially identified by Ito et al. as a deviation from the high-temperature $T$-linear behavior. There appears to be two crossover temperatures identified with the pseudogap (or spin gap) in the layered cuprates. Features in the Knight shift are identified with a temperature $T_0$ that is nearly a factor of 2 larger than the temperature $T^*$ at which the $^{63}$Cu NMR relaxation rate $1/(T_1T)$ attains a maximum. The deviation in $\rho$ (from $T$-linear behavior) is apparently closer to the higher crossover temperature $T_0$.

In principle, the opening of a pseudogap should have an observable effect on the Hall conductivity $\sigma_H$, which is sensitive to Fermi-Surface curvature. However, (as shown below) neither $\sigma_H$ nor the Hall angle $\tan \theta$ is noticeably affected at $T_0$. The most striking feature of the profile of the Hall coefficient $R_H$ (in underdoped cuprates) is its broad maximum, which occurs at a temperature much lower than $T_0$ (too low to be identified directly with $T^*$).

These puzzling inconsistencies have motivated us to reexamine the question of how the pseudogap affects the Hall effect in YBCO. Focussing on $\sigma_H$ and $\tan \theta$ rather than $R_H$, we find that the pseudogap has a strong suppressive effect on these quantities. The strong suppression first appears at a temperature coincident with $T^*$ (the peak temperature of the Cu NMR rate), rather than $T_0$. In YBa$_2$Cu$_3$O$_{6.63}$ (with $T_c \sim 60$ K), $T^* \sim 160$ K, whereas $T_0 \sim 300$ K. The Hall effect suppression extends deep into the flux-flow state. We also observe a Hall anomaly that appears in weak fields close to $T_c$. The new results bridge the interval between $T_c$ and the low-temperature flux-flow Hall results of Harris et al.

Twinned crystals of YBCO were sealed in a quartz tube with prepared polycrystalline YBa$_2$Cu$_3$O$_x$ with $x \approx 6.63$ and annealed at 500 C for 2 weeks, before quenching in liquid nitrogen ($T_c = 61.0$ and 62.5.0 K in samples 1 and 2, respectively). At each $T$, we measured both $\rho_{xx} = \rho$ and $\rho_{xy} = \rho_H$ versus a field $H \parallel c$ (with current density $J$ in-plane), and computed $\sigma_H$ and $\tan \theta = \rho_H/\rho$. 

**FIG. 1.** (Upper Panel) The temperature dependence of the in-plane Hall coefficient $R_H = \rho_H/B$ (solid symbols) in 2 crystals of twinned underdoped YBa$_2$Cu$_3$O$_{6.63}$ measured with $H \parallel c$. In both crystals, $R_H$ attains a maximum at 115 K. The open symbols are the Cu(2) NMR relaxation rate $(T_1T)^{-1}$ measured by Takigawa et al. in YBCO ($x = 6.63$). The lower panel shows $\rho$ at $H = 0$ (sample 1) and the Hall angle $\theta$ in the two crystals plotted as $T^*/\tan \theta$ with $\eta = 2.15$. Deviation of $\tan \theta$ from its high-temperature power law starts near 155 K (arrows).
In the limit \( \rho_{xx} \) and \( \rho_H \rightarrow 0 \) (the regime of interest here), \( \sigma_H \) is sensitive to alignment problems, especially near the critical temperature \( T_c \). We have adopted the following procedure to address this specific problem. At each \( T \), the field is swept from \(-8 \) T to \( 8 \) T (in 20 min.), followed by a sweep back to \( 8 \) T (the double sweeps are necessary to correct for the misalignment problems). To stabilize \( T \) to the level of \( \pm 5 \) mK (as indicated by the retracing of \( \rho \) in the four sweeps), we pre-amplified the thermometer signal before feeding it to the regulator. The measurements are then repeated at a slower sweep rate over the interval \( \pm 1 \) T to improve the resolution in weak fields. Hall curves obtained from 3 crystals are closely similar.

The Hall coefficient \( R_H \) (Fig. 1) displays a broad peak at 120 K, followed by a steep decrease as \( T \rightarrow T_c \), in good agreement with previous work \([9, 13] \). The close proximity of the maxima in \( \sigma \) suggests that the two measurements may be closely related. Instead of \( R_H \), however, it is preferable to examine the Hall angle, which provides a cleaner indicator of the Hall response.

The Cu(2) NMR relaxation rate, plotted as \( (T_1T)^{-1} \) (arb. units), suggests that the two measurements are again related. Instead of \( R_H \), however, it is preferable to examine the Hall angle, which provides a cleaner indicator of the Hall response.

In the lower panel (Fig. 3), we have plotted the resistivity \( \rho \) and the Hall angle as \( T^\eta \tan \theta \) to factor out the high-temperature dependence (the exponent \( \eta = 2.15 \) is slightly larger than the value 2.0 in 93-K YBCO \([13] \)). The plots show that \( \tan \theta \) follows the power law \( 1/T^\eta \) closely until 155 \( \pm 5 \) K where it starts deviating downwards. This temperature is quite close to the value of \( T^* \) (160 K) inferred from the NMR relaxation rate. By contrast, \( \rho \) starts deviating from its \( T \)-linear behavior at the higher temperature 250 K (we note that \( T^\eta \tan \theta \) is quite flat across this temperature.) In our crystals, the peak in \( d\rho/dT \) occurs at 175 K.

It is instructive to examine the behavior of the conductivity \( \sigma = 1/\rho \) and the (weak-field) Hall conductivity \( \sigma_H = \tan \theta/\rho \). At high \( T \), \( \sigma_H \) displays the power-law dependence \( T^{-3.85} \). Factoring out this power-law, we compare it with \( \sigma \) in Fig. 2. Whereas \( \sigma_H \) exhibits the same downturn below 155 K as the Hall angle, the \( T \)-dependence of \( \sigma_H \) remains unperturbed. These comparisons reveal that pseudogap formation affects selectively the Hall channel; they appear to have little effect on the diagonal conductivity. The selectivity again emphasizes the distinctive scattering processes that affect \( \tan \theta \) and \( \rho \) differently. Previously, the anomalous \( T \)-dependence of \( R_H \) in optimally-doped YBCO was shown to derive from distinct power-laws in \( \tan \theta \) and \( \rho \) \([13, 15] \). Here, we find that the pseudogap in underdoped YBCO selectively affects the former but not the latter.

FIG. 2. Plots of the conductivities \( \sigma_H \) (solid squares) and \( \sigma \) (open squares), with their high-\( T \) power law factored out (sample 1). The selective effect of the pseudogap on \( \sigma_H \) is apparent below 150 K. Also shown are the field scales \( H_o \) (open circles) and \( H_f \) (closed circles).

FIG. 3. (Upper Panel) The field dependence of \( R_H \) at selected \( T \) (sample 1). Below 80 K, \( R_H \) displays an anomalous decrease to zero in weak fields. The structure of this Hall anomaly evolves smoothly through the transition temperature \( T_c \). The inset shows \( \rho_H \) at 50 K at high resolution. The negative ‘dip’ that characterizes \( \rho_H \) in the flux-flow state in 93-K crystals is absent. The lower panel shows the field dependence of the Hall conductivity \( \sigma_H = \rho_H/(\rho_{xx}^2 + \rho_H^2) \) in sample 1. Below \( \sim 64 \) K, the weak-field anomaly appears as a prominent ‘missing area’ in the field profile. In the limit \( H \rightarrow 0 \), \( \sigma_H \sim H^2 \).

This selectivity also accounts for the characteristic profile of \( R_H \) in underdoped YBCO. The steep decrease in
σ_H, together with an unchanged σ (in relative terms), produces a broad peak in R_H at 120 K, followed by a rapid decrease at lower T. (We also note that the peak moves to higher T as x decreases consistent with a pseudogap effect [1].) Previously, the steep decrease in R_H below 120 K was mis-identified as the conventional Hall contribution from superconducting fluctuations [10]. The present evidence suggests that it is rather caused by the selective influence of the pseudogap on the Hall current. Fluctuations associated with superconductivity appear (with their own characteristics) only at a much lower T (80 K).

To substantiate this view, we next examine the transport behavior in both the fluctuation and flux-flow regimes. The transport features in these regimes are quite distinct from those in 93-K YBCO. In particular, we observe an anomaly in σ_H that suggests a close relation between the Hall current below T_c and in the pseudogap phase.

The upper panel of Fig. 3 shows the field dependence of R_H at selected T. At and above 80 K, R_H is independent of field, within our resolution. Below 80 K, however, a weak-field anomaly becomes apparent as a rapid decrease in R_H to zero in the limit H → 0. At temperatures above T_c, the steep variation of R_H in weak fields terminates quite abruptly at a characteristic field H_n above which R_H is H-independent. At 70 K, for example, we estimate H_n ≥ 0.6 T (arrow) (H_n is plotted vs. T in Fig. 2). The Hall anomaly becomes quite striking below T_c. In the lower panel, we have plotted the field dependence of σ_H. While the anomaly is barely resolved in weak fields at 70 K, it grows to a prominent ‘missing area’ in the field profile below 60 K. Both above and below T_c, σ_H has the limiting form sgn(H)H^n (as H → 0), with an anomalous exponent n = 2 ± 0.2.

![Graph of ρ vs. μ_H](image)

**FIG. 4.** (Upper Panel) The magnetoresistance in sample 1 near 80 K shown at high resolution (sample 1). The vertical scales are indicated by the double arrows.

The rather abrupt appearance of the Hall anomaly at 80 K is matched by the appearance of a weak-field anomaly in the MR (Fig. 3). At and above 83 K, ρ displays only the weak H^2 MR associated with the normal state. Below 83 K, however, a weak-field cusp is observed. It is worth a remark that the cusp (and the weak-field Hall anomaly) represent the first field-sensitive transport features that become observable as T is lowered from 300 K. Below T_c, the cusp magnitude increases rapidly as the resistivity evolves into the usual flux-flow profile. The onset of dissipation in the resistive state is marked by a fairly well-defined threshold field H_f (the curve for H_f is displayed in Fig. 3).

Both the Hall anomaly and the cusp in the MR are clearly associated with the field suppression of superconducting droplets as we approach T_c. However, we emphasize that these two features are poorly described by the standard fluctuation conductivity calculations based on Aslamasov-Larkin (AL) theory. Whereas the AL theory explains reasonably well the fluctuating conductivity measured in 93-K crystals, the field dependence of ρ here is far too large and rapid to be described by these perturbation results. The fairly abrupt appearance of the weak-field cusp at 80 K is quite different from the powerlaw dependence (on reduced temperature ε) of the AL fluctuation terms [10].

At temperatures above T_c, σ_H is also inconsistent with the conventional fluctuation picture. The weak-field exponent n is incompatible with the fluctuation Hall conductivity calculated, for example, from time-dependent Ginzburg Landau (GL) theory, which has the H-linear form [10]

\[ \delta \sigma_H = \left( \frac{e^2}{3\pi d} \gamma_n \right) \frac{H \xi_0^2}{h e^2}, \]

where \( \xi_0 \) is the GL coherence length, and \( \gamma_n \) and g are GL parameters.

More significantly, the Hall current in underdoped YBCO in its mixed state above 40 K is strikingly different from that in 93-K YBCO (over the same range in reduced T). It is well known that, in 93-K YBCO, the flux-flow Hall resistivity exhibits a sign change that appears abruptly below T_c as a prominent negative ‘dip’ in the profile of ρ_H vs. H [15,13,19]. The non-monotonic ρ_H simplifies considerably when converted to σ_H [18,19]; the latter is simply the sum of a positive quasiparticle (qp) term \( \sigma_H^{qp} \) and a negative vortex-flow term \( \sigma_H^f \), viz.

\[ \sigma_H = \sigma_H^{qp} + \sigma_H^f, \]

with H dependences

\[ \sigma_H^{qp} \sim H, \quad \sigma_H^f \sim -1/H. \]  \hspace{1cm} (1)

The \( 1/H \) divergence in \( \sigma_H^f \) is a firm prediction of many theories [20]. Experimentally, it shows up in 93-K YBCO as the characteristic negative ‘dip’ in \( \rho_H \).

By contrast, \( \rho_H \) in 60-K YBCO shows no evidence of the negative ‘dip’ feature above 40 K. At 50 K, for instance, the onset of dissipation (ρ) occurs abruptly at 0.35 T. However, the Hall resistivity \( \rho_H \) is unresolved from zero until about 0.8 T, above which it increases in the positive direction (inset in lower panel of Fig. 3). The absence of the \( 1/H \) divergence in \( \sigma_H \) is also apparent in the traces in the lower panel of Fig. 3. Other aspects of \( \sigma_H \) are also different from that in Eq. 1. Instead of increasing monotonically, \( \sigma_H \) is nominally \( H^2 \) at low fields.
and attains a broad maximum at high fields. Thus, between 40 and 60 K, the vortex Hall conductivity term \( \sigma_H \) is conspicuously absent. This is possibly the most striking difference observed so far in vortex transport properties between underdoped and optimally-doped YBCO. It implies that, in the former above 40 K, the vortices move very nearly parallel to \( \mathbf{J} \times \mathbf{H} \), i.e. the flux-flow electric field contributes only to \( \rho_{xx} \). (Only below 30 K does the vortex Hall contribution become apparent \([3]\). Below 20 K, the vortex Hall angle rapidly approaches the superclean regime.)

The absence of \( \sigma_H \) at the temperatures where the Hall anomaly is most prominent persuades us that the latter reflects changes to the electronic state, as reflected in \( \sigma_{\text{qp}} \). We note that, on both sides of \( T_c \) (Fig. 3), \( \sigma_H \) exhibits the same anomalous \( H^2 \) dependence in low fields. In our view, the continuous behavior of the Hall anomaly across \( T_c \) implies that it represents the same feature in the electronic current on both sides of \( T_c \). The traces in Fig. 3 imply that the increased phase coherence near \( T_c \) leads to an electronic state (below the line \( H_n \)) in which the Hall conductivity is severely suppressed. The similarity of this suppression to the Hall-angle suppression at 155 K is highly suggestive. The high-temperature suppression appears to be a less complete precursor of the more complete suppression that occurs below 80 K (but it is more robust in field).

In conclusion, we find that the pseudogap onset temperature \( T_\ast \), as derived from \( 1/(T_\ast T)^{\ast} \), correlates closely with a strong downturn in \( \sigma_H \) and \( \tan \theta \), whereas \( \rho \) is relatively unaffected. Within the broad interval 80 to 160 K, in which the selective effect of the pseudogap on the two currents is fully apparent, the effects of superconducting fluctuations are either non-existent or too small to be resolved by transport experiments. Fluctuations effects become apparent only below 80 K, where the Hall anomaly and the cusp in the MR first appear. At the high temperature side, the fluctuation regime is clearly demarcated by the field \( H_n \). The rapid growth of phase coherence produces a weak-field anomaly in \( \sigma_H \) that corresponds to a further suppression of the Hall conductivity. In the mixed state, our measurements show that the contribution of the vortices to the Hall current is observable above 40 K. This implies that the suppression of \( \sigma_H \) is associated with the quasiparticle current. Unlike in 93-K crystals, the pseudogap state displays several unusual transport features that evolve systematically between 160 K and 40 K. These results seem to provide a more complete picture of the effect of the pseudogap on the transport properties in underdoped YBCO.

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