High-Field Quasiparticle Tunneling in Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$: Negative Magnetoresistance in the Superconducting State

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We report on the c-axis resistivity $\rho_c(H)$ in Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ that peaks in quasi-static magnetic fields up to 60 T. By suppressing the Josephson part of the two-channel (Cooper pair/quasiparticle) conductivity $\sigma_q(H)$, we find that the negative slope of $\rho_c(H)$ above the peak is due to quasiparticle tunneling conductivity $\sigma_q(H)$ across the CuO$_2$ layers below $H_{c2}$. At high fields, the field at the peak in $\rho_c(H)$ is controlled by QP tunneling in the superconducting state of Bi-2212. We access QP current by suppressing the Josephson current in two ways: (i) with transport current in thin Bi-2212 mesas, and (ii) with a 60 T field, which in underdoped Bi-2212 crystals exposes QP tunneling down to $\sim 22$ K. We show that high-field c-axis conductivity is linear-in-field, in agreement with the field-linear QP conductivity $\sigma_q$ we find in Bi-2212 mesas in the resistive state. Near 60 T, at low temperatures $\rho_c$ tends to a finite (saturation) value, as predicted for a d-wave superconductor. The crossover to the normal state above $T_c$ is witnessed by a superlinear $\sigma_q(H)$, plausibly related to the pseudogap.

A clue to the mystery of high temperature superconductivity is likely to arrive from the most peculiar normal-state properties of cuprate superconductors [1]. The observed ‘pseudogap’ features [2] in the quasiparticle (QP) spectrum of the underdoped cuprates, such as Bi$_2$Sr$_2$CaCu$_2$O$_{8+\delta}$ (Bi-2212) [3-6] and YBa$_2$Cu$_3$O$_{7-\delta}$ [7], are taken by many as a signature of a non-Fermi-liquid electronic structure above $T_c$ [7]. On the other hand, recent c-axis tunneling [8] data in the superconducting state of Bi-2212 suggest that the d-wave Fermi-liquid character of QPs is restored at low temperatures and low magnetic fields. The question arises where a crossover to the non-Fermi liquid behavior will occur and what is its signature in the c-axis transport – a direct probe of QP tunneling. Energies of the order of the gap $\Delta_0$ can be accessed with temperature, applied voltage, or magnetic field. The effect of field may be unique, since it can affect other degrees of freedom (e.g., magnetic excitations) that are coupled to QPs.

The field dependence of the c-axis resistivity was first considered by Briceño et al. and by Gray and Kim [1] in the context of a ‘giant’ magnetoresistance vs temperature peak observed in Bi-2212. They ascribed the $\rho_c(T)$ peak to two competing processes: (1) Josephson and (2) quasiparticle tunneling, the latter masked by superconductivity at the nodes of the $d$-wave Fermi-surface [9]. The former, understood only recently by Koshelev [10], leads to a positive c-axis magnetoresistance at sufficiently low fields [11,12]. At high fields, the observed crossover to a negative slope in $\rho_c(H)$ has been attributed to the normal state [11,14] and to the emergence of the pseudo- or spin-gap [13]. Consequently, the field at the peak in $\rho_c(H)$ in Bi-2212 was taken as ‘$H_{c2}$’ [14], in an apparent conflict with the $H$-dependencies of $\rho_{ab}(H)$ and in-plane $T_c(H)$ [11,12].

In this Letter we resolve the origin of negative c-axis magnetoresistance in layered high-$T_c$ superconductors by exploring quasiparticle dissipation in magnetic fields up to 60 Tesla. We demonstrate that at high fields, above the maximum in $\rho_c(H)$ at $H^*$, negative magnetoresistance in Bi-2212 is controlled by QP tunneling in the superconducting state. We access QP current by suppressing the Josephson current in two ways: (i) with transport current in thin Bi-2212 mesas, and (ii) with a 60 T field, which in underdoped Bi-2212 crystals exposes QP tunneling down to $\sim 22$ K. We show that high-field c-axis conductivity is linear-in-field, in agreement with the field-linear QP conductivity $\sigma_q$ we find in Bi-2212 mesas in the resistive state. Near 60 T, at low temperatures $\rho_c$ tends to a finite (saturation) value, as predicted for a d-wave superconductor. The crossover to the normal state above $T_c$ is witnessed by a superlinear $\sigma_q(H)$, plausibly related to the pseudogap.

We used high quality, optimally and slightly underdoped Bi-2212 crystals ($T_c$’s $\approx 92.5$ and 89 K respectively). Here we will show data on the underdoped crystal (900 $\times$ 600 $\times$ 20 $\mu$m$^3$), since it afforded us a downward expanded temperature range of negative magnetoresistance. The resistivity was measured in magnetic fields along the c-axis up to 60 T generated in the Long Pulse (LP) System at the National High Magnetic Field Laboratory (NHMFL) in Los Alamos, NM. One advantage of the LP system is that it allows us to avoid sample heating by the induced eddy currents – the main problem encountered in short-pulse experiments, where a large $dB/dt \geq 10^4$ T/s produces significant heating and associated thermal hysteresis effects for larger than micron-size samples. The LP system delivers a nearly triangular field pulse with the fastest up- and down-ramp rates $\sim 150$ and 360 T/s and a duration $\sim 2$ sec. A non-hysteretic resistivity at these different rates assures a minimal effect of eddy currents. The temperature was controlled to better than 50 mK, with the massive copper-dust/epoxy sample holder serving as a thermal anchor. A standard 4-probe contact configuration was used to measure resistance by a lock-in technique at 17 kHz, recording it with a fast analog-to-digital converter. The current-voltage ($I$-
V) characteristics of mesa-patterned high quality Bi-2212 whiskers, prepared by a Focused Ion Beam technique [8], were measured at NHMFL in Tallahassee, FL.

Figure 1 shows the raw data of the c-axis resistivity $\rho_c$ vs magnetic field for the Bi-2212 crystal in the temperature range 22.5 K - 110 K. We point to two major features below $T_c = 89$ K: first, a maximum in $\rho_c(H)$ and second, the near saturation of high-field $\rho_c$ vs $T$ at the lowest $T$ (inset in Fig. 1) [15]. The latter is a new observation. Note that above 35 K, $\rho_c(T)$ at 55 T roughly follows a $\ln T$ dependence up to $\sim T_c$. Above $T_c$, $\rho_c$ has a weaker ($\sim 0.1\% / T$ at 110 K) field dependence. Remarkably, a 60 T field even at 70 K does not suppress $\rho_c$ to it’s value above $T_c$.

![Figure 1](image_url)

**FIG. 1.** Out-of-plane resistivity $\rho_c$ vs $H$ of Bi-2212 crystal in fields up to 60 T at different temperatures. $j = 0.05$ A/cm². Inset: $\rho_c$ vs $T$ at 55 T (●) and at 0 T (line).

A maximum in $\rho_c(H)$ at $H^*(T)$ is summarized in the $H - T$ diagram in Fig. 2. Below $H^*$ the magnetoresistance is positive and above negative. Such $H$-dependence has been observed before [12][14], and attributed to a crossover to the normal state. Indeed, since there is a (weak) negative magnetoresistance above $T_c$ (Fig. 1), one may easily fall into the trap of identifying $H_{c2}$ with $H^*$. However, $\rho_c$ is a measure of tunneling across the CuO$_2$ layers and a direct signature of the normal state should come from the in-plane (non-tunneling) resistivity $\rho_{ab}(H)$. It reveals a maximum [12] at $H > H^*$ (inset in Fig. 2): at $H^*$, $\rho_{ab}$ has still a strong positive magnetoresistance, clearly originating from the superconducting state. So, this is a first telltale sign that $H_{c2} > H^*$, since the changes in $\rho_{ab}(H)$ and $\rho_c(H)$ do not track. Thus we take the maximum in $\rho_{ab}(H)$ as an underestimate of $H_{c2}$ [10] (see Fig 2), the assignment we will confirm in the considerations that follow.

Below we show that a key to understanding magnetoresistance is the realization that the c-axis conductivity is a parallel, two-channel tunneling process: (i) $\sigma_J$ of Cooper pairs (mostly at low fields) and (ii) $\sigma_q$ of quasiparticles (dominant at higher fields): $\sigma_c = \sigma_J + \sigma_q$. We first discuss $\sigma_J$. Consider a stack of superconducting sheets spaced by $s$ subject to a c-axis magnetic field. Dissipation for Josephson interlayer current is caused by phase difference slips due to motion of the Josephson strings attached to pancake vortices in adjacent layers [10]. Diffusive drift of pancakes governed by pinning and thermal fluctuations results in motion of current-driven Josephson strings. The barrier for pancake motion is large at low fields but lessens with increasing field [3]. The field dependence is explicit in the derived universal relationship between the in-plane and the c-axis conductivity [10]

$$\sigma_J \propto \sigma_{ab} \Phi_0 s^2 E_J^2 / (BT^2) \propto \sigma_{ab} / B,$$

where $E_J$ is the Josephson energy per unit area at $B = 0$, $B \sim H$, and $\sigma_{ab} \propto \exp(U/T)$ is controlled by the energy barriers $U(H)$ to thermally activated pancake hopping [16]. For a 2D vortex lattice $U(H) \sim -\ln H$ [16], implying a power-law field dependence of $\sigma_J = a H^{-\gamma}$ with $\gamma(T) \approx 1 + U/T$. The low-field $\sigma_{ab}/\sigma_c$ is indeed nearly linear in $H$ (inset in Fig. 3), confirming Eq. 1.

![Figure 2](image_url)

**FIG. 2.** $H - T$ diagram showing $H^*(T)$, the location of peak in $\rho_c(H)$ (●). Rough estimates of $H_{c2}(T)$ from the broad maximum in $\rho_{ab}(H)$ [12] (▲), from the demise of the $j$-dependence of $\rho_c(H)$ (●, see Fig. 3), and from $1/\beta$ (▼) in Eq. 2. Inset: $\rho_c$ and $\rho_{ab}$ vs $H$. A positive $d\rho_{ab}/dH$ persists well above $H^*(T)$.

The Josephson nature of $\sigma_J$ implies that this contribution and thus $H^*$ may be significantly affected by an injection of extra interlayer transport current $j$. This current – in addition to pancake disorder – suppresses Josephson coupling, while its effect on $\sigma_q$ is negligible. For the data of Fig. 1, a downshift of $H^*$ at larger $j$ is evident in Fig. 3. At a higher temperature (55 K) the $j$-dependence of $\rho_c(H)$ (i.e., Josephson current) smoothly disappears above 45 T. This crossover into an ohmic dis-
sipation regime – controlled mainly by the QPs – further marks a lower bound to $H_d > H^*$ (Fig. 2).

An immediate consequence of this picture is that at high fields $\tau_f \ll \tau_q$, and thus $\tau_s \sim \tau_q$. To unambiguously separate the field behavior of the two channels, we first independently test the field dependence of $\sigma_q$ via $c$-axis $I$-$V$ characteristics in a tiny ($\sim 2 \mu m^2$) step-like mesa bar (sketch in Fig. 4) carved out of a Bi-2212 whisker, with the thickness of $\sim 50$ CuO$_2$ layers. Here, the Josephson current can be suppressed by a fairly low transport current and the decreasing branch of the $I$-$V$ curves in the resistive state is determined purely by quasiparticle tunneling. The inset in Fig. 4 shows a set of $I$-$V$’s from which QP conductivity (the initial slope of $\sigma$) was extracted as in Ref. [8]. The obtained $\sigma_q(H,T)$ vs $c$-axis DC field is remarkably linear up to 33 T at all temperatures (main panel of Fig. 4).

We turn now to the entire field and temperature dependence of the $c$-axis conductivity in a macroscopic crystal

$$\sigma_c(H,T) \approx \alpha H^{-\nu} + \sigma_q(0,0)[\eta + \beta H].$$

Fig. 5 shows $\sigma_c(H)$ for the crystal. Below $T_c$ there is clearly a decreasing $\tau_f$ below $H^*$ and, for $H \gg H^*$, the $H$-linear $\sigma_q$. Both channels are well described by Eq. 2 as illustrated by a fit at 55 K. Such fits give $\nu$ in the range 1.5 - 3.5 between 70 and 22.5 K, confirming the $\ln H$ dependence of hopping barriers with $U \sim 45$ K at low $T$. The coefficient $\alpha(T)$ increases with decreasing $T$ [14]. The temperature dependence of the zero-field $\sigma_q(0,T)$ gives $\eta(T) = 1 + c T^2$ (inset in Fig. 5), in agreement with the results on mesas. From the coefficient $c = \pi^2/18\gamma^2 (T < \gamma)$ [8] we find the effective scattering rate of QPs due to impurities inside layers, $\gamma \sim 0.4 T_c$.

And, significantly, we obtain a nonzero extrapolated $\sigma_q(H \to 0, T \to 0) \approx 2.5 \text{ (k}\Omega \text{ cm)}^{-1}$. This result is a signature of a $d$-wave superconductor [20] and it confirms direct measurements of $\sigma_q(0,0) \sim 1.5 - 3.7 \text{ (k}\Omega \text{ cm)}^{-1}$ in mesas [6]. This explains the saturation of high-field $\rho_c$ at $T \to 0$ (Fig. 1) which naturally reflects the nonzero $\sigma_q$ arising from the scattering by impurities in the nodal regions of the gap. Above $T_c$, the two-channel description of Eq. 2 obviously fails. The $\sigma_q(H)$ becomes weaker and is no longer linear (Fig. 5). The change across $T_c$ is very gradual and subtle, consistent with the existence of the pseudogap above $T_c$ [20].

We now address our result – the linear increase of $\sigma_q$ with $H$. Within the Fermi-liquid model, $\sigma_q(H)$ in the superconducting state is proportional to the QP density of states (DOS) and inversely proportional to the effective scattering rate for interlayer tunneling, $\gamma_c$. In the $s$-wave picture, $\sigma_q \approx \sigma_N H/H_{c2} (\sigma_N$ is the normal state conductivity), since the DOS is proportional to the number of vortex cores. However, in a $d$-wave superconductor the dominant contribution to $\sigma_q$ comes from the fourfold-symmetry nodal protrusions of delocalized states outside the cores [21]. They also lead to the increase of DOS with $H$ due to the Doppler-shifted QP energy caused by the in-plane supercurrents around vortices. Recent calculation [22] shows that leading order in the Doppler shift energy with respect to $\Delta_0$, the increase of DOS is compensated by an increase of $\gamma_c$. The increase of $\gamma_c$ is caused by random positions of pancake vortices resulting in different Doppler shifts between equivalent points in adjacent layers. This compensation results in a field independent $\sigma_q(H,T) \approx \sigma_q(0,0) H/T$ in Eq. 2. But, at high fields the core contributions and other corrections (i.e., nonlinear QP spectrum near gap nodes) [23] must enter. These corrections give a net increase of $\sigma_q(H)$

FIG. 3. Current dependence of $\rho_c(H)$ at 55 K and 35 K. At a higher current density $\rho_c$ reaches maximum at a lower field. Inset: The ratio $\sigma_{ab}/\sigma_c$ vs $H$ at low fields, as in Eq.[1].

FIG. 4. Normalized quasiparticle $c$-axis conductivity as a function of $H \parallel c$ obtained from the $I$-$V$ curves (top inset) measured on the mesa-shaped Bi-2212 (sketched).
as $\sim \sigma_q(0,0)H/H_\Delta$. $H_\Delta = \Phi_0\Delta^2_f/hv_F^2$ is $\approx 40-80$ T at low $T$. In the BCS theory $H_\Delta$ corresponds to $H_{c2} \sim \Phi_0/(\pi\xi^2)$ and $1/\beta$ (Eq. 2) gives a rough estimate of $H_{c2}(T)$ plotted in Fig. 2. However, this estimate is fuzzy since the above does not include the pseudogap which may complicate the magnetoresistance at high fields.

Note that at a 60 T field $\sigma_q(H,T)$ is only $\approx 2\sigma_q(0,0)$, far below the normal state value: e.g., at 140 K $\sigma_c \approx 55(\text{k}\Omega\text{ cm})^{-1} \sim 20\sigma_q(0,0)$. It is also below $\sigma_N$ reached by applying voltage $V > 2\Delta_0/e$ at low $T$. This is surprising. It may reflect the influence of the pseudogap which has been shown to be robust to changes in the magnetic field. This is consistent with the gapped QP spectrum inside vortex cores observed by STM.

In summary, the field dependence of the $c$-axis conductivity in Bi-2212 is consistently described by tunneling of Cooper pairs and of quasiparticles across a $d$-wave superconductor. We demonstrate that the maximum in $\rho_c(H)$ comes from the competition between these two conduction channels below $H_{c2}$. The low-$T$ saturation of the high-field $\rho_c(T)$ is a consequence of the $d$-wave character of the superconducting state. Remarkably, $\sigma_c(H)$ remains linear in $H$ and small up to 60 T. It becomes gradually weaker and superlinear above $T_c$, where the QP tunneling is controlled by the pseudogap.

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\*FIG. 5. $\sigma_c(H)$ below and above $T_c = 89$ K. A fit at 55 K to a superposition of Cooper pair (dash) and quasiparticle (dash-dot) contributions in Eq. 2 is indicated. The high-field $\sigma_c(H)$ becomes nonlinear above $T_c$. Inset: Zero-field $\sigma_q$ vs $T/T_c$ extracted from $\sigma_c(H,T)$ for a Bi-2212 crystal for $j = 0.05$ A/cm$^2$ (●) and $j = 0.1$ A/cm$^2$ (⊙), and obtained from the $I$-$V$ curves for a mesa (◦). Both fit a $T^2$-dependence (dash) up to $T \sim \gamma$.\*