High-Field Quasiparticle Tunneling in $Bi_2Sr_2CaCu_2O_{8+\delta}$: Negative Magnetoresistance in the Superconducting State

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We report on the c-axis resistivity $\rho_c(H)$ in $Bi_2Sr_2CaCu_2O_{8+\delta}$ that peaks in quasistatic magnetic fields up to 60 T. By suppressing the Josephson part of the two-channel (Cooper pair/quasiparticle) conductivity $\sigma_c(H)$, 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 (a) $\sigma_q(H)$ grows linearly with H, and (b) $\rho_c(T)$ tends to saturate ($\sigma_c \neq 0$) as $T \rightarrow 0$, consistent with the scattering at the nodes of the d-wave gap. A superlinear $\sigma_q(H)$ marks the normal state above T_c.

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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 $Bi_2Sr_2CaCu_2O_{8+\delta}$ (Bi-2212) [3,4] and $YBa_2Cu_3O_{7-\delta}$ [5,6], 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 Δ_0 can be accessed with temperature, applied voltage, or magnetic field. The effect of the 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 [9] 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 below T_c. The former, understood only recently by Koshelev [10], leads to a positive c-axis magnetoresistance at sufficiently low fields [9,11,12]. 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 [12]. 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* $T_{c2}(H)$ [9,11,12].

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 T. We demonstrate that at high fields, above the maximum in $\rho_c(H)$ at $H^\star$, negative magnetoresistance in Bi-2212 is controlled by QP tunneling in the superconducting state. We access the 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 σ_q we find in Bi-2212 mesas in the resistive state. Near 60 T, at low temperatures ρ_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 ($900 \times 600 \times 22\ \mu m^3$), 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 of $\sim$150 and 360 T/s and a duration of $\sim$2 s. A nonhysteretic 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 four-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

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ion beam technique [8], were measured at NHMFL in Tallahassee, FL.

Figure 1 shows the raw data of the c-axis resistivity ρ_c vs magnetic field for the Bi-2212 crystal in the temperature range 22.5–110 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 ρ_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, ρ_c has a weaker ($\sim 0.1\%/\mathrm{T}$ at 110 K) field dependence. Remarkably, a 60 T field even at 70 K does not suppress ρ_c to its value above T_c.

A maximum in $\rho_c(H)$ at $H^\star(T)$ is summarized in the H-T diagram in Fig. 2. Below $H^\star$ 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 $H^\star$. However, ρ_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 (nontunneling) resistivity $\rho_{ab}(H)$. It reveals a maximum [12] at $H > H^\star$ (inset in Fig. 2): at $H^\star$, ρ_{ab} has still a strong positive magnetoresistance, clearly originating from the superconducting state. So, this is a first telltale sign that $H_{c2} > H^\star$, 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} [16] (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) σ_J of Cooper pairs (mostly at low fields) and (ii) σ_q of quasiparticles (dominant at higher fields): $\sigma_c = \sigma_J + \sigma_q$. We first discuss σ_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 the motion of the Josephson strings attached to pancake vortices in adjacent layers [10]. A diffusive drift of pancakes governed by pinning and thermal fluctuations results in a motion of current-driven Josephson strings. The barrier for pancake motion is large at low fields but lessens with increasing field [9]. The field dependence is explicit in the derived universal relationship between the in-plane and the c-axis conductivity [10],

$$\sigma_J \propto \sigma_{ab}\Phi_0 s^2 E_J^2/(BT^2) \propto \sigma_{ab}/B\,, \tag{1}$$

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 [17]. For a 2D vortex lattice $U(H) \sim -\ln H$ [18], implying a power-law field dependence of $\sigma_J = \alpha H^{-\nu}$ with $\nu(T) \approx 1 + U/T$. The low-field σ_{ab}/σ_c is indeed nearly linear in H (inset in Fig. 3), confirming Eq. (1).

The Josephson nature of σ_J implies that this contribution and thus $H^\star$ 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 σ_q is negligible. For the data of Fig. 1, a down-shift of $H^\star$ 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 dissipation regime, controlled mainly by the QPs [19], further marks a lower bound to $H_{c2} > H^\star$ (Fig. 2).

An immediate consequence of this picture is that at high fields $\sigma_J \ll \sigma_q$, and thus $\sigma_c \sim \sigma_q$. To unambiguosly

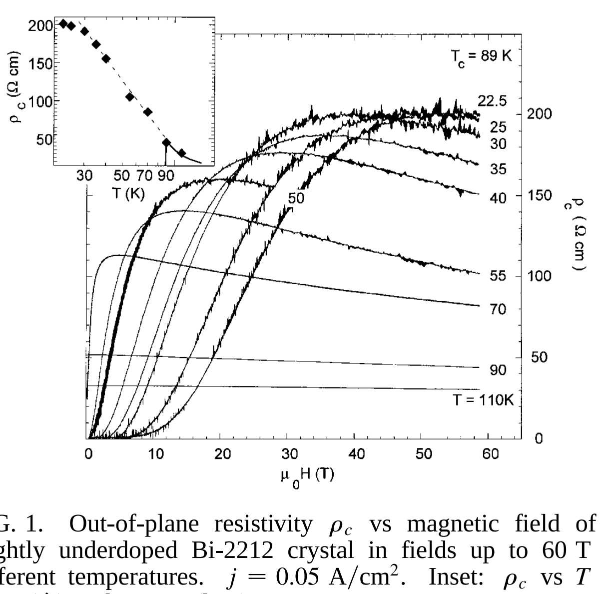

FIG. 1. Out-of-plane resistivity ρ_c vs magnetic field of a slightly underdoped Bi-2212 crystal in fields up to 60 T at different temperatures. $j = 0.05$ A/cm^2. Inset: ρ_c vs T at 55 T (◆) and at 0 T (line).

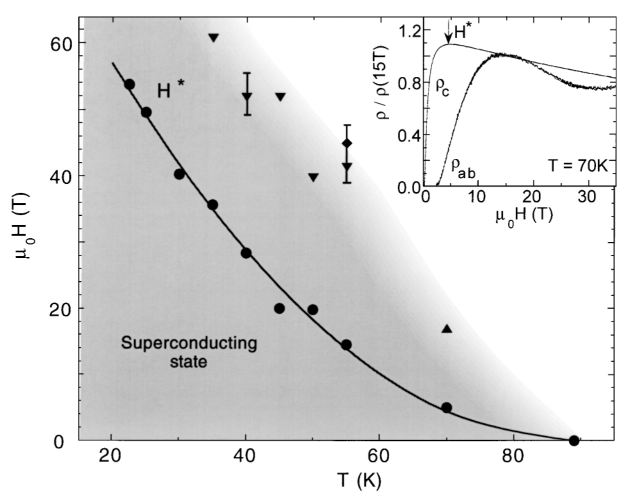

FIG. 2. H-T diagram showing $H^*(T)$, the location of the peak in $\rho_c(H)$ (•). A lower bound to $H_{c2}(T)$ is from the broad maximum in $\rho_{ab}(H)$ [12] (▲) and from the demise of the j dependence of $\rho_c(H)$ in Fig. 3 (◆). $1/\beta$ (▼) from the fits to Eq. (2) can be taken as a rough estimate of $H_{c2}(T)$ (see text). Inset: ρ_c and ρ_{ab} vs H at 70 K. A *positive* $d\rho_{ab}/dH$ persists well above $H^*(T)$.

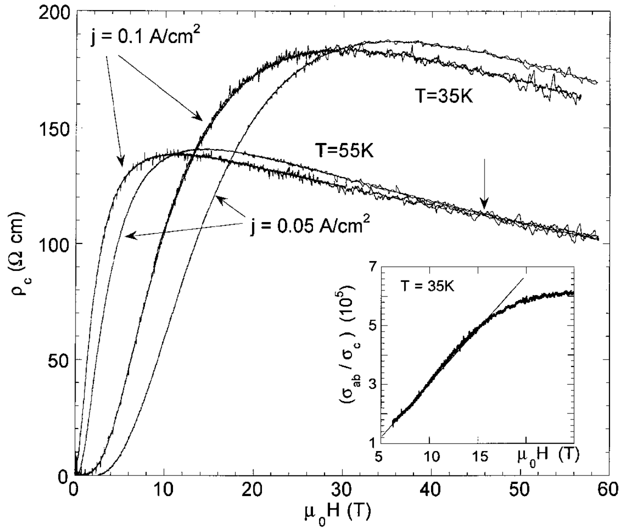

FIG. 3. Current dependence of $\rho_c(H)$ at 55 and 35 K. At a higher current density ρ_c reaches the maximum at a lower field. At 55 K, the j dependence vanishes $\gtrsim$45 T (arrow). Inset: The ratio σ_{ab}/σ_c vs H at low fields, as in Eq. (1).

separate the field behavior of the two channels, we first *independently* test the field dependence of σ_q via c-axis I-V characteristics in a tiny ($\sim$2 μm^2) steplike 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 [8] 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 I vs V) was extracted as in Ref. [8]. The obtained $\sigma_q(H,T)$ vs the 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,

$$\sigma_c(H,T) \approx \underbrace{\alpha H^{-\nu}}_{\sigma_J} + \underbrace{\sigma_q(0,0)[\eta + \beta H]}_{\sigma_q}. \qquad (2)$$

Figure 5 shows $\sigma_c(H)$ for the crystal. Below T_c there is clearly a decreasing σ_J below $H^\star$ and, for $H \gg H^\star$ the H-linear σ_q. Both channels are well described by Eq. (2) as illustrated by a fit at 55 K. Such fits give ν 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 [17]. The temperature dependence of the zero-field $\sigma_q(0,T)$ gives $\eta(T) = 1 + cT^2$ at low T (inset in Fig. 5), in agreement with the results on mesas. From the coefficient $c = \pi^2/18\gamma^2$ ($T \ll \gamma$) [8] we find the effective scattering rate of QPs due to impurities inside layers, $\gamma \sim 0.4T_c$. Significantly, we obtain a *nonzero* extrapolated $\sigma_q(H \to 0, T \to 0)$ [$\approx$2.5 (kΩ cm)$^{-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 (kΩ cm)$^{-1}$ in mesas

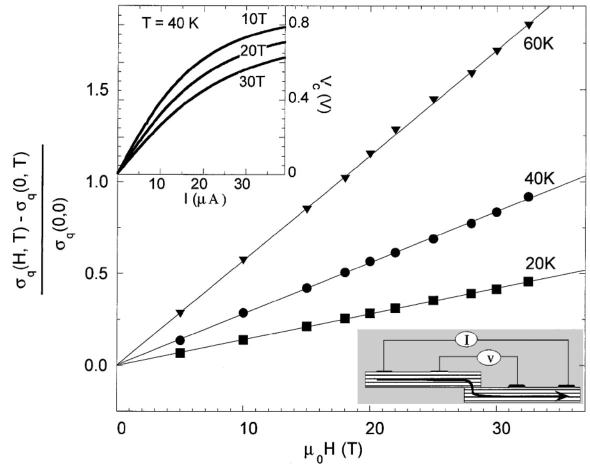

FIG. 4. Normalized quasiparticle c-axis conductivity as a function of $H \parallel c$ obtained from the I-V curves (top inset) measured on the mesa-shaped Bi-2212 (sketched).

[8]. This explains the saturation of high-field ρ_c at $T \to 0$ (Fig. 1) which naturally reflects the nonzero σ_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 [2,3].

We now address our result—the linear increase of σ_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, γ_c. In the

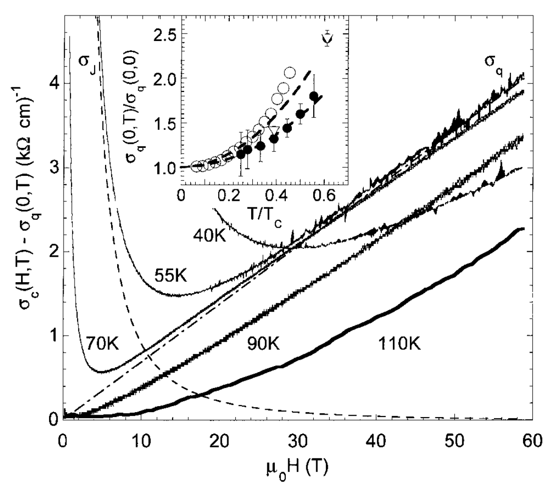

FIG. 5. $\sigma_c(H)$ below and above $T_c = 89$ K. A fit at 55 K to a superposition of Cooper pair (dashed line) and quasiparticle (dash-dotted line) contributions in Eq. (2) is indicated. The high-field $\sigma_c(H)$ becomes nonlinear above T_c ($\sim H^{1.2}$ at 90 K and $\sim H^{1.6}$ at 110 K). Inset: Zero-field σ_q vs T/T_c extracted from $\sigma_c(H,T)$ for a Bi-2212 crystal for $j = 0.05$ A/cm^2 ($\bullet$) and $j = 0.1$ A/cm^2 (∇), and obtained from the I-V curves for a mesa ($\circ$). Both fit a T^2 dependence (dashed line) up to $T \sim \gamma$.

s-wave picture, $\sigma_q \approx \sigma_N H/H_{c2}$ (σ_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 σ_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. A recent calculation [22] shows that, to leading order in the Doppler shift energy with respect to Δ_0, the increase of DOS is compensated by an increase of γ_c. The increase of γ_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 $\sigma_q(H,T) \approx \sigma_q(0,0)\eta(T)$ in Eq. (2). But, at high fields the core contributions and other corrections (e.g., nonlinear QP spectrum near gap nodes in [22]) must enter. These contributions give a net increase of $\sigma_q(H)$ as $\sim\sigma_q(0,0)H/H_\Delta$. $H_\Delta = \Phi_0\Delta_0^2/\hbar^2 v_F^2$ is $\approx$40–80 T at low T [22]. In the BCS theory H_Δ corresponds to $H_{c2} \sim \Phi_0/2\pi\xi^2$ and $1/\beta$ [Eq. (2)] gives a rough but consistent estimate of $H_{c2}(T)$ plotted in Fig. 2. Note that this estimate is fuzzy, since the above considerations do not include the pseudogap which may complicate the magnetoresistance at high fields.

Surprisingly, at a 60 T field $\sigma_q(H,T)$ is only $\approx 2\sigma_q(0,T)$, far below the normal state value: e.g., at 140 K, $\sigma_c \approx 55(\mathrm{k\Omega\ cm})^{-1} \sim 20\sigma_q(0,0)$. It is also below σ_N reached by applying voltage $V > 2\Delta_0/e$ at low T [8,23]. It may reflect the influence of the pseudogap which has been shown to be robust to changes in the magnetic field [5]. This is consistent with the gapped QP spectrum inside vortex cores observed by scanning tunneling microscopy [3].

In summary, the field dependence of the c-axis conductivity in Bi-2212 is consistently described by the 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, $\sigma_q(H)$ 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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