Unconventionally large quantum-dissipative gap regime in overdoped $Bi_2Sr_2CaCu_2O_{8+y}$

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We report on an unexpectedly extensive dissipative gapped regime in the zero-temperature (quantum) limit in highly overdoped $Bi_2Sr_2CaCu_2O_{8+y}$. Using high-field interlayer resistivity measurements we show that as $T\to 0$ the magnetic field that closes the pseudogap and the upper critical field H_{c2} coincide, uniquely defining the upper limit on the vortex state. By mapping the upper and lower bounds on the molten vortex state, we find the gapped quantum fluctuation regime stretching from ~ 30 to 70 T. This exceeds by far the conventional estimates, pointing to the anomalous gapped nature of the strongly overdoped regime.

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High-temperature superconductivity in cuprates arises from doping the charge carriers into a Mott insulating parent compound. So, naturally, the spotlight is chiefly trained on the low doping (underdoped) regime where the pseudogap is the feature most prominent and the least understood. The pseudogap—an incomplete depletion in the quasiparticle density-of-state (DOS) near the Fermi energy[1]—dominates the phase diagram way above the superconducting transition temperature T_c, expressing striking electronic correlations preserving the *d*-symmetry of the superconducting gap below T_c but with finite dissipation. By a way of contrast, on the high doping (overdoped) side cuprates are deemed conventional, and as such are expected to display behaviors of conventional (Fermi-liquid) metals.

The very root of high T_c is intertwined with the existence of the pseudogap viewed in disparate theoretical proposals either as a form of a predecessor[2–4] or as a competitor[5,6] to high T_c. So, on the one hand, the pseudogap onset temperature $T^\star(>T_c)$ may characterize the onset of some sort of pairing correlations[2,3], while the overall phase coherence is set at T_c. Indeed, vortexlike excitations in cuprates were reported[7] well above T_c. On the other hand, the "competitor" scenarios require the pseudogap to be closed in a phase transition at the so-called "quantum critical point" (QCP),[5,6] turning the cuprate pseudogapless on the overdoped side.[8] Thus, central to the competing views is the very existence of the pseudogap in strongly overdoped superconducting cuprates—the existence experimentally still in question, since with increased doping the pseudogap temperature $T^\star$ rapidly falls and when it arrives to the neighborhood of T_c the pseudogap is arguably hard to distinguish from the usual thermal fluctuation effects.[9]

Here we focus instead on the dissipative gapped state (DGS) in the zero temperature (quantum) limit that can be accessed at ultrahigh fields. In the quantum regime, the superconducting state sets in below the upper critical field H_{c2}, whereas the superconducting coherence is globally established when the dissipationless state is reached. This occurs when "vortex matter" solidifies[10–12]—conventionally very near H_{c2} (Refs. 10 and 13). Evaluating the significance of quantum fluctuations is a harder task, and estimating (very high) H_{c2} in the cuprates has been a subject of much controversy,[14] not surprisingly aggravated by the pseudogap below $T^\star$. We adopt an approach based on closing the pseudogap with the field H_{pg} and restoring the superconductor to its ungapped "true" normal state, since this sets an "ultimate" upper limit[7] to the vortex state. Probing the gap with the intrinsic (between CuO_2 planes) tunneling *c*-axis resistivity[15] we track the field H_{sc} at which tunneling of quasiparticles overtakes that of Cooper pairs.[16] This field sets a lower bound on H_{c2}. We show that in highly overdoped $Bi_2Sr_2CaCu_2O_{8+y}$ (BSCCO) these lower and upper bounds merge at $T=0$, unambiguously fixing the value of $H_{c2}(0)$. By mapping the zero-resistivity field $H_{0\rho}$ we find the DGS near $T=0$ (the regime of quantum vortex melt) even far on the overdoped side *unconventionally large*, unaccountable by the usual phenomenology of the mixed state in the quantum limit.

We chose a BSCCO crystal with the hole doping $p=0.225$ and—essential for sizing up the fluctuation effects—a sharp diamagnetic transition at $T_c(0)\approx 60$ K (Fig. 1). To overdope, the crystal was annealed to high homogeneity in 200 atm O_2. Measurements were performed at 100 kHz in a 60 ms pulse magnet at National High Magnetic Field Laboratory (NHMFL) in Los Alamos using a lock-in technique. A negligible eddy-current heating was verified by the consistency of data taken with successive pulses with different maximum target fields.

The crystal displays the expected sharp "second peak" magnetization anomaly[17,18] nearly up to T_c at a field $H_{sp}\lesssim 0.1$ T (Fig. 1) that marks a transition from a low-field linelike [three-dimensional (3D)] ordered vortex lattice to a high-field pancakelike [two-dimensional (2D)] disordered solid.[19] Our point of departure is the 2D vortex solid above H_{sp} which driven by fluctuations[10] will melt into a liquid distinguished by a finite dissipation. This crystal is so overdoped that the zero-field *c*-axis resistivity $\rho_c(T)$ is metallic $(d\rho_c(T)/dT>0)$ nearly all the way down to T_c (Fig. 2). However, the application of even a moderate magnetic field

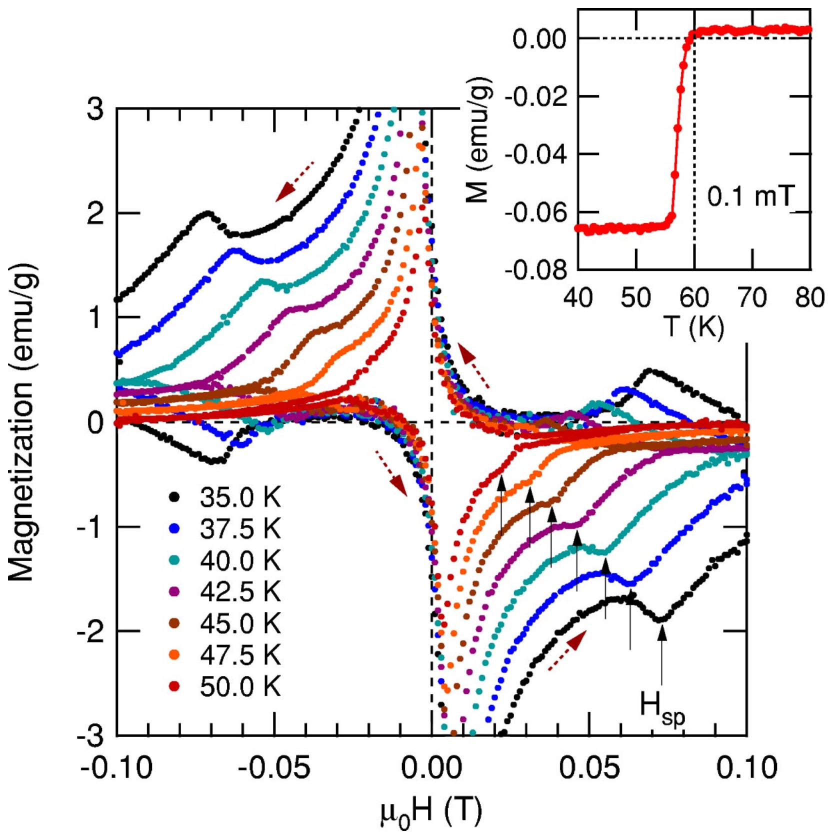

FIG. 1. (Color online) Magnetization M vs H loops in a highly overdoped BSCCO crystal at low fields (dotted arrows indicate the field-sweep direction). A characteristic "second peak" is at $H_{sp}(T)$ (arrows). In contrast with optimally doped BSCCO (Ref. 17), in such overdoped regime this peak associated with an order-disorder transition is always in the irreversible regime (Ref. 18). The sharp second peak and the sharp magnetic transition at T_c (inset) ensure high crystal quality and a homogeneous doping.

($\sim$10 T) along the c direction uncovers the semiconducting upturn in $\rho_c(T)$ before it plunges to zero in the dissipationless state below T_c. An upward deviation of $\rho_c(T)$ from the metallic $\rho_c^n(T)$ is a measure of the DOS depletion at low excitation energies, since the intrinsic tunneling conductivity $\sigma_c(E)$ is a measure of the DOS in the CuO_2 planes. Further increases in the field begin to affect the pseudogap itself; the upturn is suppressed and the metallic regime is extended to lower temperatures. The metallic regime in the fully ungapped state is characterized by the temperature-linear $\rho_c^n(T)$

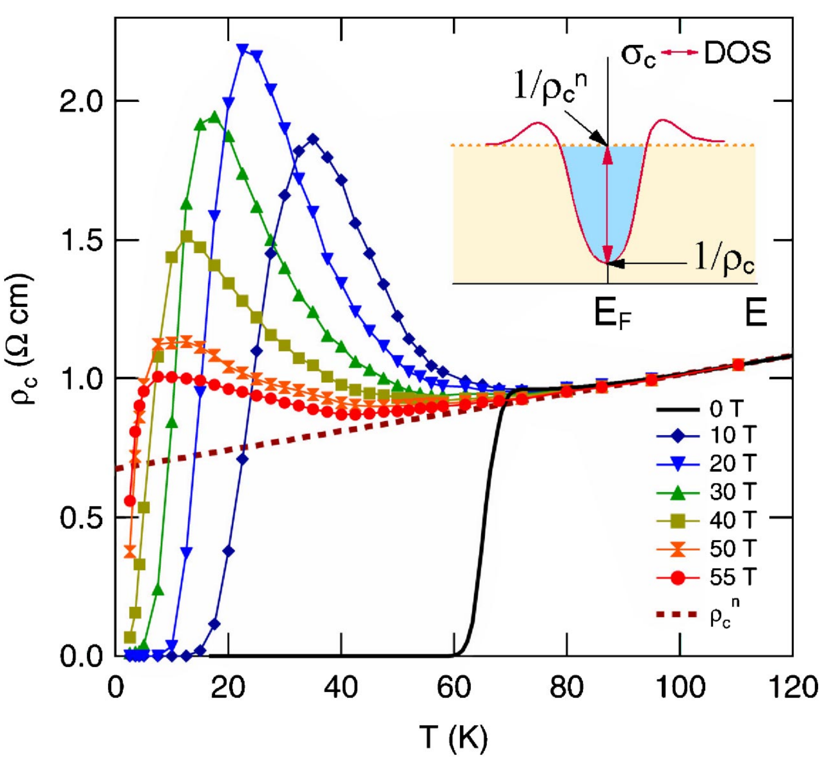

FIG. 2. (Color online) c-axis resistivity vs temperature in overdoped BSCCO up to 55 T. The normal ungapped state resistivity $\rho_c^n(T)$ (dashed line) is evaluated from the measured T-linear dependence (Ref. 20) above $T^\star$ ($\sim$100 K). Inset: ρ_c corresponds to the inverse of the interlayer tunneling (differential) conductivity σ_c near Fermi energy $E=E_F$.

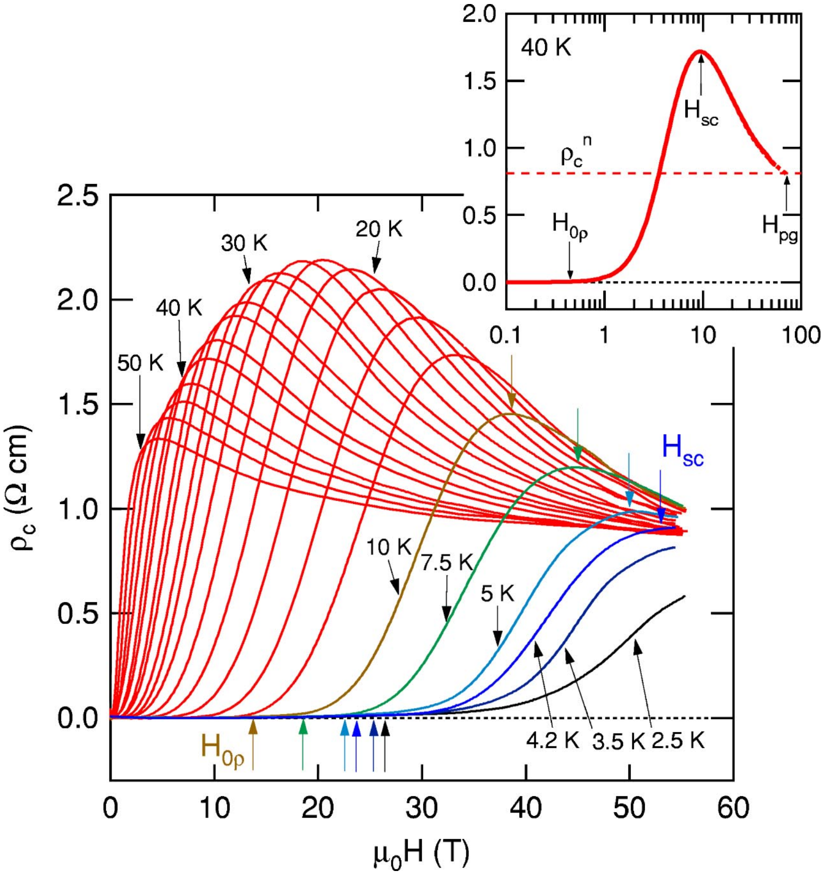

FIG. 3. (Color online) c-axis resistivity $\rho_c(H)$ is marked by three characteristic fields: $H_{0\rho}$, H_{sc}, and H_{pg}. The onset of finite ρ_c at $H_{0\rho}$ is a harbinger of a vortex liquid state. Above the peak field $H_{sc}(\lesssim H_{c2})$, quasiparticle tunneling becomes dominant. The negative magnetoresistance (MR) above $H_{sc}(T)$ disappears at the pseudogap temperature $T^\star$ (Ref. 15). ρ_c reaches its ungapped state value ρ_c^n (dashed line) at the pseudogap closing field H_{pg} (inset, see text). The $\rho_c(H)$ curves broaden at low T, more so below 10 K. Large quantum fluctuations appear as a large ''gap'' between $H_{0\rho}$ and H_{sc}.

(Fig. 2) down to very low temperatures.[20] The excess quasiparticle resistivity $\Delta\rho_c$ obtained by subtracting the ρ_c^n was recently shown to follow a power-law field dependence up to 60 T.[15] Upon closing the pseudogap $\Delta\rho_c \rightarrow 0$, and we determine the pseudogap closing field $H_{pg}(T)$ by comparing $\rho_c(H)$ and ρ_c^n at each temperature through the power-law fit, as described in Ref. 15.

To probe the DGS from ''below,'' we turn to lower fields where ρ_c is zero (Fig. 3), the vortex state is a solid with finite pinning, and our crystal is truly superconducting. An earmark of the solid-to-liquid transformation in CuO_2 planes is the onset of finite dissipation due to mobile two-dimensional (pancake) vortices.[14] The pancakes in adjacent planes are connected by Josephson strings that are pulled along, generating phase slips responsible for the dissipation in the c direction.[16] ρ_c sensitively detects the vortex agitation in the planes and we use it to evaluate the onset of the liquid phase at $H_{0\rho}(T)$. For the estimate of $H_{0\rho}$ we took a $\rho_c = 0.01\rho_c^n$ criterion[21] appropriate to our experimental resolution, see the inset in Fig. 3. With a further increase in field, $\rho_c(H)$ reaches a maximum at a temperature-dependent field $H_{sc}(T)$, where the Josephson (Cooper-pair tunneling) current becomes comparable to the tunneling of quasiparticles.[16] $H_{sc}(T)$ can be accurately tracked defining a limit on the H_{c2} ($\gtrsim H_{sc}$) from below.

Let us now dissect the complete field-temperature diagram shown in Fig. 4(a) by the practiced phenomenology of ''vortex matter'' physics. The upper critical field must lie somewhat above H_{sc} but below H_{pg}. However, as $T \rightarrow 0$, H_{sc} and H_{pg} merge, giving a firm evaluation of $H_{c2}(0) \approx 70$ T. It is evident that the low-temperature values of $H_{0\rho}$ are well below $H_{c2}(0)$—-a nearly 40 T region between these two limits is the regime of quantum vortex liquid. In forming the liquid (and producing a finite ρ_c) quantum fluctuations have to overcome (at least) the elastic energies in the vortex system. The quantum regime is identified by the relevant energy $\hbar\omega_0$ of the quantum vibration modes which can be estimated as follows. Taking the vortex dynamics to be dissipative[22] with viscosity $\eta \approx (H_{c2}B)/\rho_n c^2$ (induction $B \sim H$ at high fields, ρ_n is the in-plane dissipation in the normal state, and c is the velocity of light), we compare elastic $[\propto c_{66}(4\pi/a_0^2)]$ and dynamic $[\propto \eta\omega]$ energies to find a typical $\omega_0 \approx \rho_n c^2 B/(16\pi\lambda^2 H_{c2})$. Here, $c_{66} = \Phi_0 B/(8\pi\lambda)^2$ is the shear modulus, $a_0 = (\Phi_0/B)^{1/2}$ is the intervortex spacing, and λ is the magnetic penetration depth. Close to $H_{c2}(0)$, with $\lambda \sim 200$ nm and $\rho_n \sim 20\ \mu\Omega$cm (Ref. 23) we obtain $\hbar\omega_0 \sim 10$ K. So at high fields and below 10 K, quantum effects dominate the melting process.

The general Lindemann criterion says that the solid turns unstable when the mean vortex displacement $\langle u^2 \rangle^{1/2} = c_L a_0$ is a good fraction of a_0 with $c_L \approx 0.1$–0.3 (Ref. 14). For $k_B T > \hbar\omega_0$, we have the equipartition theorem $\langle u^2 \rangle \approx 0.2T/(c_{66}d)$, where d is the interlayer spacing. And as near H_{c2} the superfluid density is suppressed, the thermal melting line of a 2D solid becomes

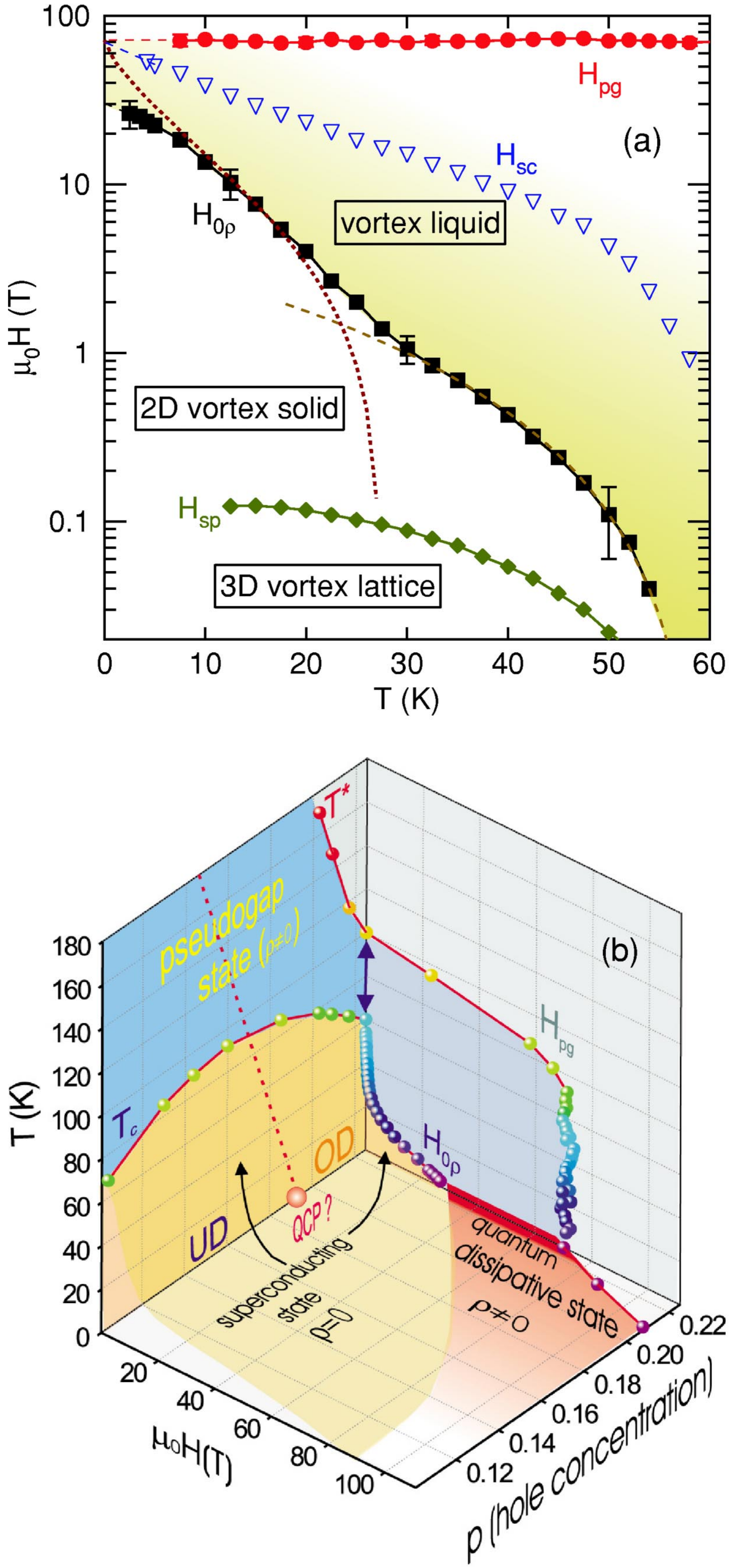

FIG. 4. (Color online) (a) H-T diagram for overdoped BSCCO reveals a huge regime of quantum melt (shaded) with finite dissipation. The previously unexplored regime is above the 3D-2D transition (Ref. 19) at $H_{sp}(T)$. As $T\rightarrow0$, $H_{sc}(T)$ is exponential in temperature (Ref. 15) with no saturation, approaching ~70 T—the same as the T-independent H_{pg} (error bars in $H_{pg}(T)$ are the size of symbols). This fixes the value of $H_{c2}(0)$, exposing quantum melt within $[H_{c2}(0),\ H_{0\rho}]$. Nearer T_c thermal melting expectedly follows the $(1-T/T_c)^2$ dependence (Ref. 14) (dashed line). At ~25 K, $H_{0\rho}(T)$ crosses to the 2D thermal melting in Eq. (1) (dotted line), with $H_{c2}(T)$ approximated by $H_{sc}(T)$. (b) Temperature-field-doping diagram of BSCCO highlights the overdoped (OD) region near T_c and near $T=0$. In the scale shown, at $H=0$ the dissipative pseudogapped state (DGS) between $T^{\star}$ "line" and T_c fills the T-p phase space. At $T=0$ the DGS between $H_{0\rho}$ and H_{pg} occupies unconventionally large region in the sample more overdoped than the proposed QCP near $p=0.19$.

$$T_m^{th}/T_c \approx (c_L^2/Gi^{2D})(1-H/H_{c2})^2, \tag{1}$$

where the Ginzburg number Gi^{2D} (with $d\sim 15$ Å, $Gi^{2D}\sim 0.05$, see Ref. 14) is a measure of thermal fluctuations in two dimensions. This expression with $c_L=0.14$ above 10 K accounts for the low-temperature trend of $H_{0\rho}(T)$.

At $T=0$, a rough Lindemann estimate in the 2D limit[10] gives quantum melting at

$$H_m^Q/H_{c2} \approx 1-0.6\sqrt{\frac{\omega_{max}}{\omega_0}}\exp\left(\frac{-\pi^3 c_L^2 R_Q}{4R_\square}\right), \tag{2}$$

where $R_Q=\hbar/e^2=4.1$ kΩ is the (universal) quantum resistance, $R_\square=\rho_n/d\sim 150\ \Omega$, and $\omega_{max}=\Delta/\hbar$ is set by the superconducting gap $\Delta\simeq 20$ meV (Ref. 1). The observed over $\sim$50% suppression of $H_{0\rho}$ as $T\to 0$ corresponding to $H_m^Q/H_{c2}\lesssim 0.5$ can only be accounted for with a very small $c_L\lesssim 0.1$.

Such a small Lindemann number (or such large size of the quantum melt) is unexpected in the context of models requiring the strongly overdoped regime to be conventional. It is the classic 2D dislocation mediated (Kosterlitz-Thouless) thermal melting that corresponds to $c_L\approx 0.1$. Quantum melting of a 2D lattice should be harder,[24] requiring $c_L\approx 0.25$— a number closely obtained, e.g., in layered (2D) organic superconductors[25,26] that also have a d symmetry.[27] Note that, following Ref. 28, disorder pushes the line of zero resistivity (quantum glass transition) towards higher field values. With our rough criterion, $H_{0\rho}$ in Fig. 4 significantly overestimates the irreversibility field,[21] reducing the value of c_L even further. The observed ease of vortex displacements by zero-point vibrations naturally points to a modified structure of the vortex core. The pseudogapped core has been experimentally demonstrated by scanning tunneling spectroscopy.[29] And from phenomenological considerations, the observed reduction in the effective viscosity η (Ref. 30) [implying a higher vortex velocity and thus a larger $R_\square$ in Eq. (2)], points to a reduced number of carriers available for pushing the current through the core, consistent with the pseudogap depletion in the DOS.

The T-p-H diagram is compiled in Fig. 4(b), where we focus on a regime not accessed previously by other techniques. In a view where the deduced QCP is near $p=0.19$,[8] the difference between T_c and $T^\star$ (double-ended arrow) beyond the QCP may come from thermal fluctuations. However, the unconventionally large DGS we observe at high fields as $T\to 0$ may suggest the pseudogap in the quantum limit far on the overdoped side. This picture is consistent with the upturn in the $\rho_c(T)$ recovered in a moderate magnetic field, and the smooth continuity of $H_{pg}(p)$ from the underdoped (UD) side[15] with its ''flat'' low-temperature behavior at *all* doping levels.

In summary, by exploring the phase space below the pseudogap closing field we obtain a reliable value of the upper critical field H_{c2} in the quantum limit. A clear quantitative consequence of this is that the estimates that work well for the conventional superconductors fail to account for the observed large size of the gapped quantum fluctuation regime, witnessing an anomalous behavior of a strongly overdoped cuprate.

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