Two-stage magnetic-field-tuned superconductor–insulator transition in underdoped La$_{2-x}$Sr$_x$CuO$_4$

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In the underdoped pseudogap regime of cuprate superconductors, the normal state is commonly probed by applying a magnetic field ($H$). However, the nature of the $H$-induced resistive state has been the subject of a long-term debate, and clear evidence for a zero-temperature $H$-tuned superconductor–insulator transition has proved elusive. Here we report magnetoresistance measurements on underdoped La$_{2-x}$Sr$_x$CuO$_4$, providing striking evidence for quantum-critical behaviour of the resistivity—the signature of a $H$-driven superconductor–insulator transition. The transition is not direct, being accompanied by the emergence of an intermediate state, which is a superconductor only at temperature $T = 0$. Our finding of a two-stage $H$-driven superconductor–insulator transition goes beyond the conventional scenario in which a single quantum critical point separates the superconductor and the insulator in the presence of a perpendicular magnetic field. Similar two-stage $H$-driven superconductor–insulator transitions, in which both disorder and quantum phase fluctuations play an important role, may also be expected in other copper-oxide high-temperature superconductors.

The superconductor–insulator transition (SIT) is an example of a quantum phase transition (QPT): a continuous phase transition that occurs at $T = 0$, controlled by some parameter of the Hamiltonian of the system, such as doping or the external magnetic field. A QPT can affect the behaviour of the system up to surprisingly high temperatures. In fact, many unusual properties of various strongly correlated materials have been attributed to the proximity of quantum critical points (QCPs). An experimental signature of a QPT at nonzero $T$ is the observation of scaling behaviour with relevant parameters in describing the data. Although the SIT has been studied extensively, even in conventional superconductors many questions remain about the perpendicular-field-driven SIT in two-dimensional (2D) or quasi-2D systems. In high-$T_c$ cuprates (here $T_c$ is the transition temperature), which have a quasi-2D nature, early magnetoresistance experiments showed the suppression of superconductivity with high $H$, revealing the insulating behaviour and hinting at an underlying $H$-field-driven SIT (ref. 7). However, even though understanding the effects of $H$ is believed to be essential to understanding high-$T_c$ cuprate superconductivity and continues to be a subject of intense research, the evidence for the $H$-field-driven SIT and the associated QPT scaling in cuprates remains scant and inconclusive.

In the conventional picture of type-II superconductors, $H$ penetrates the sample in the form of a solid lattice of interacting vortex lines in the entire mixed state $H_1(T) < H < H_2(T)$, where $H_1$ is the Meissner field and $H_2$ is the upper critical field. This picture, however, neglects fluctuations which, in high-$T_c$ superconductors, are especially important. Indeed, the delicate interplay of thermal fluctuations, quantum fluctuations and disorder leads to a complex $H$–$T$ phase diagram of vortex matter. Theoretical calculations, for example, suggest melting of the vortex solid into a vortex liquid for fields below what is now a crossover line $H_c(T)$. Quantum fluctuations could result in a vortex liquid persisting down to $T = 0$. At very low $T$, the disorder becomes important and modifies the vortex phase diagram such that there are two distinct vortex solid phases: a Bragg glass with $T_c(H) > 0$ at lower fields and a vortex glass with $T_c = 0$ at higher fields. The SIT would then correspond to a transition from a vortex glass to an insulator at even higher $H$. However, the interplay of this vortex line physics and quantum criticality, the key question in the $H$-field-driven SIT, has remained largely unexplored.

In this study, we find strong evidence for the $H$-field-driven SIT in underdoped La$_{2-x}$Sr$_x$CuO$_4$ (LSCO). The results are consistent with the existence of three phases at $T = 0$, although the behaviour of the in-plane resistivity $\rho$ suggests the presence of the direct SIT over a surprisingly wide range of $T$ and $H$. At the lowest $T$, however, $\rho(T, H)$ data reveal an intermediate phase with $T_c = 0$ and the true SIT at higher $H$. We focus on samples with a relatively low $T_c(H = 0)$ to ensure that experimentally attainable $H$ are high enough to fully suppress superconductivity. Unlike most other studies, ours includes samples grown using different techniques (Methods), to separate out any effects that may depend on the sample preparation conditions from the more general behaviour. As well as providing evidence for the SIT, the magnetoresistance data are used to calculate the contribution of superconducting fluctuations (SCFs) to conductivity, allowing a construction of the $H$–$T$ phase diagram.

In-plane resistivity of La$_{2-x}$Sr$_x$CuO$_4$

One sample was a film with the nominal composition La$_{1.93}$Sr$_{0.07}$CuO$_4$ (ref. 11) and a measured $T_c(H = 0)$ of ($3.8 \pm 0.1$) K. The high-quality single crystal$^1$ had $T_c(H = 0) = (5.2 \pm 0.1)$ K and the nominal composition La$_{1.9}$Sr$_{0.1}$CuO$_{4+y}$, with $y$ not precisely known (see Methods for more details). Unless otherwise specified, $T_c$ is defined throughout as the temperature at which the resistance (that is, $\rho$) becomes zero for a given $H$. The method to determine $T_c(H)$ is illustrated in Supplementary Fig. 1.)

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Figure 1 | Temperature dependence of the in-plane resistivity $\rho$ in different magnetic fields $H$ for $x = 0.07$ LSCO film. a, The $\rho(T,H)$ data exhibit a change of the sign of $\rho/T$ as a function of $H$ at high $T \geq 5$K ($R_{\parallel}/\rho_{\parallel}$ is resistance per square per CuO$_2$ layer; see Methods). Except for $H=0$, the solid lines are guides for the eye. b, A selection of the $\rho(T,H)$ data from a, shown on a log–log scale to focus on the low-$T$, intermediate-$H$ regime. Short-dashed lines guide the eye. Solid lines represent power-law fits $\rho(T,H) = \rho_0(H) T^{\alpha(H)}$. c, Fitting parameters $\rho_0(H)$ and $\alpha(H)$. Error bars indicate 1 standard deviation (s.d.) in the fits for $\rho_0(H)$ and $\alpha(H)$. Short-dashed lines guide the eye.

Figure 1a shows the $\rho(T)$ curves for different $H$ that were extracted from the magnetoresistance measurements at fixed $T$ for the $x = 0.07$ sample (see Supplementary Fig. 2 for similar data on the $x = 0.06$ crystal). A small $H$ clearly leads to a decrease of $T_c(H)$, such that $T_c \rightarrow 0$ for $H \approx 4$T (Supplementary Information and Supplementary Fig. 3). The survival of the superconducting phase with $T_c > 0$ up to fields much higher than the Meissner field $H_M \approx 100$ Oe is understood to be a consequence of the pinning of vortices by disorder[10]. As $H$ increases further, a pronounced maximum appears in $\rho(T)$. The temperature at the maximum, $T_{\text{max}}$, shifts to lower $T$ with increasing $H$, similar to early observations[1]. At the highest $H$, $\rho(T)$ curves exhibit insulating behaviour.

The intermediate $H$ regime, in which each $\rho(T)$ curve exhibits a maximum, is especially intriguing. Here the system shows a tendency towards insulating behaviour already at high $T > 5$K, but then the sign of $\rho/T$ changes, suggesting that another mechanism sets in at lower $T$ and drives $\rho(T \rightarrow 0)$ towards zero. The low-$T$, intermediate-$H$ regime is more evident in Fig. 1b, where the same data are presented on a log–log scale. For $T < T_{\text{max}}$, the data are described best with the phenomenological power-law fits $\rho(T,H) = \rho_0(H) T^{\alpha(H)}$, which indicate that $\rho(T = 0) = 0$—that is, that the system is a superconductor only at $T = 0$. The exponent $\alpha$ depends on $H$ (Fig. 1c) and goes to zero at $H \sim 13.5$T. This implies the existence of a $T$-independent $\rho$ at that field. The non-mono-tonic behaviour of $\rho(T)$ in the intermediate regime suggests that high-$T$ $(T > T_{\text{max}})$ and low-$T$ $(T < T_{\text{max}})$ regions should be examined separately and more closely.

Scaling analysis of the high-temperature behaviour

Figure 2a shows a more detailed set of magnetoresistance measurements on the $x = 0.07$ film, focused on the behaviour at $5$K $\leq T < 10$K. The magnetoresistance curves clearly exhibit a well-defined, $T$-independent crossing point at $\mu_s H_c^* = 3.63$T. In other words, this is the field where $\partial \rho/\partial T$ changes sign from positive or metallic at low $H$, to negative or insulating at $H > H_c^*$ (see also Fig. 1a). We find that, near $H_c^*$, $\rho(T)$ for different $H$ can be collapsed onto a single function by rescaling the temperature, as shown in Fig. 2b. Therefore, $\rho(T,H) = \rho(T/H_c^*)$, where the scaling parameter $T/H_c^*$ is found to be the same function $\delta = (H - H_c^*)/H_c^*$ on both sides of $H_c^*$. In particular, $T_c \propto |\delta|^{1/2}$ with $zv \approx 0.73$ (Fig. 2c) over a remarkably wide, more than two orders of magnitude range of $|\delta|$. Such a single-parameter scaling of the resistance is precisely what is expected$^{11}$ near a $T = 0$ SIT in 2D, where $z$ and $v$ are the dynamical and correlation length exponents, respectively$^{2}$. Similar scaling on the $x = 0.06$ crystal sample (Supplementary Fig. 4) yields a comparable $zv \approx 0.59 \pm 0.08$. It is interesting that, although the critical fields $H_c^*$ in the two samples differ by almost a factor of two ($\mu_s H_c^* = 3.63$T for $x = 0.07$ film and $\mu_s H_c^* = 6.68$T for $x = 0.06$ crystal), the critical resistivities $\rho_c^*$ are almost the same ($R_{\parallel}$, low $\approx (17.4 \pm 18.0) \Omega$; see also Supplementary Information). The exponent product $zv \sim 0.7$ has been observed for perpendicular $H$-field-tuned transitions also in conventional 2D superconductors (for example, in a-Bi (ref. 14) and a-NbS$_2$ (ref. 15)) and, more recently, in 2D superconducting LaTiO$_3$/SrTiO$_3$ interfaces$^{16}$. The value $zv \sim 0.7$ is believed to be in
the universality class of the 2D SIT in the clean limit, as described by the (2 + 1)D XY model and assuming that z = 1 owing to the long-range Coulomb interaction between charges13,17.

Scaling analysis of the low-temperature behaviour

As shown above, the behaviour of the system over a range of T above T_{max} seems to be controlled by the QCP corresponding to the transition from a superconductor to an insulator in the absence of disorder, and driven by quantum phase fluctuations. As T is lowered below T_{max}, however, this transition does not actually take place, as some of the insulating curves assume the power-law dependence ρ(T, H) = ρ_n(T)T^{α(0)} (Fig. 1), leading to a superconducting state at T = 0. Figure 3a shows a set of magnetoresistance measurements carried out at very low T < T_{max}, which exhibit a T-independent crossing point at µ_0H^*_c = 3.63 T, consistent with α(H = H^*_c) ≈ 0 (see also Fig. 1b,c). Near H^*_c, an excellent scaling of ρ with temperature according to ρ(T, H) = ρ_n(T/T^*) is obtained (Fig. 3b), where T^*_c ∝ |δ|^z on both sides of H^*_c (Fig. 3c). Here δ = (H - H^*_c)/H^*_c and zv = 1.15 ± 0.05. Assuming z = 1, this type of single-parameter scaling with ν = 1 corresponds to the T = 0 SIT in a 2D disordered system19.

Superconducting fluctuations

The extent of SCFs can be determined from the transverse (H || c) magnetoresistance by mapping out fields H^*_c(T) above which the normal state is fully restored11,18-20. In the normal state at low fields, the magnetoresistance increases as H^2 (ref. 21), so that the values of H^*_c can be found from the downward deviations from such quadratic dependence that arise from superconductivity when H < H^*_c. The H^*_c(T) line determined using this method (Supplementary Fig. 5) is shown in Fig. 4a for the x = 0.07 film15 (see Supplementary Information and Supplementary Fig. 6 for the x = 0.06 crystal sample). In both cases, H^*_c(T) is well fitted by H^*_c(T) = H^*_c(0)[1 - (T/T^*)^2], where the normal-state resistivity ρ_n(T, H) was obtained by extrapolating the region of H^2 magnetoresistance observed at high enough H and T, as illustrated in Supplementary Fig. 5.

Phase diagram

In addition to H^*_c(T) and Δρ_{SCF}(T, H), the phase diagram in Fig. 4a includes T_c(H) and T_{max}(H), as well as the critical fields H^*_c and H^*_c. For both samples, T_c(H = H_0) = 0 (Supplementary Information and Supplementary Fig. 3) for fields H_0 ∝ H^*_c (µ_0H^*_c = 3.63 T for x = 0.07 film and µ_0H^*_c = 6.68 T for x = 0.06 crystal). This is consistent with the T = 0 transition from a pinned vortex solid to another phase at higher fields. For the x = 0.07 film sample, T_{max}(H = H_0) = 0 for µ_0H^*_c = (13.4 ± 0.7) T, meaning that T_{max} vanishes at the critical field µ_0H^*_c = 13.45 T within the measurement error. In the x = 0.06 single crystal, T_{max}(H) extrapolates to zero at µ_0H^*_c = (14.0 ± 0.6) T— that is, at a field similar to that in the film sample, even though their H^*_c(T = 0) are very different. The vanishing of a characteristic energy scale, such as T_{max}(H), in the T = 0 limit is consistent with the existence of a quantum phase transition at H^*_c.
The scaling regions associated with the critical fields \( H^*_1 \) and \( H^*_2 \) are also shown in Fig. 4a. It is striking that the ‘hidden’ critical point at \( H^*_2 \), corresponding to the SIT in the clean limit, dominates a huge part of the phase diagram. (See also Supplementary Fig. 7 for the QCP \( H^*_1 \).) Two QCPs, \( H^*_1 \) and \( H^*_2 \), correspond to the two QCPs in Fig. 4b, which includes crossover temperatures \( T^*_c \) fitted to a VRH form, but the change of the behaviour at the lowest \( T^*_c \) on the other hand, is remarkably small (Fig. 4a), but it dominates and attributed to 2D or 3D Mott VRH. 

The scaling region corresponding to the critical point at \( H^*_2 \), on the other hand, is remarkably small (Fig. 4a), but it dominates the behaviour at the lowest \( T^*_c \). On the insulating side—that is, for \( H > H^*_2 \)—there is evidence for the presence of SCFs, as the applied \( H \) increases. The lowest-\( T^*_c \) insulating behaviour (Fig. 3b) may also be fitted to a VRH form, but the change of \( \rho \) is too small again to determine the VRH exponent with high certainty.

A simplified version of the same phase diagram is shown in Fig. 4b, which includes crossover temperatures \( T^*_1 \) and \( T^*_2 \), corresponding to the two QCPs \( H^*_1 \) and \( H^*_2 \), respectively. The logarithmic \( T \)-scale also makes it more apparent that the vortex solid phase with \( T^*_c(H) \) extending precisely up to \( H^*_1 \)—that is, that the QCP \( H^*_1 \) is associated with the quantum melting of the pinned vortex solid. The other phases shown in Fig. 4b are discussed below.

### Discussion

To account for our observation of three distinct phases as \( T \to 0 \) and two QPTs, we discuss our results in the context of other relevant work on the same material (see Fig. 5 for a sketch of the phase diagram, which is based mainly on Fig. 4 and Supplementary Fig. 6). In particular, we note that two order parameters are needed for a complete description of the behaviour of a type-II superconductor in the presence of \( H^*_1 \). One that describes superconductivity and another one that describes vortex matter (see, for example, ref. 23).

At low fields below the \( T^*_c(H) \) line, \( \rho(T) = 0 \), which is attributed to the pinning of the vortex solid\(^{10-12} \). \( T^*_c(H) \) is known\(^{20,25} \) to correspond to the so-called irreversibility temperature \( T_{\text{irr}}(H) \) below which the magnetization becomes hysteretic, indicating a transition of the vortex system between a low-\( T \), low-\( H \) pinned regime and an unpinned one\(^{26} \). This low-\( H \) vortex solid phase has been identified experimentally as a Bragg glass in LSCO (ref. 26), as well as in other cuprates\(^{8,10} \) and some conventional superconductors (for example, 2\( H \)-NbSe\(_2\) (ref. 27)). The Bragg glass forms when the disorder is weak\(^{28,29} \); it retains the topological order of the Abrikosov vortex lattice (see sketch of a Bragg glass in Fig. 5) but yields broadened diffraction peaks. As such a distorted Abrikosov lattice has many metastable states and barriers to motion, it is, strictly speaking, a glass.

At higher \( H \), where the density of vortices is larger\(^{28,29} \), a topologically disordered, amorphous vortex glass is expected (see sketch of a vortex glass in Fig. 5). A transition from a Bragg glass to a vortex...
Figure 4 | Transport H–T phase diagram and scaling regions in underdoped LSCO. The data are shown for \( x = 0.07 \) film. a, The colour map (on a log scale) shows the contribution of superconducting fluctuations (SCF) to conductivity \( \Delta \sigma_{SCF} \) versus \( T \) and \( H \). The dashed red line is a fit with \( \mu_0 H_c(T) = 15(T - T(0))/29^2 \). The error bars indicate the uncertainty in \( H_c \) that corresponds to 1 s.d. in the slopes of the linear fits in Supplementary Fig. 5. The horizontal dashed black lines mark the values of the \( T = 0 \) critical fields \( H_0 \) and \( H_2 \) for scaling. The pink lines show the high-\( T \) scaling region: the hashed symbols mark areas beyond which scaling fails, and the dots indicate areas beyond which the data are not available. The green lines show the low-\( T \) scaling region. b, Simplified phase diagram showing the crossover temperatures \( T_1^* \) and \( T_2^* \) corresponding to \( H_0 \) and \( H_2 \), respectively, and the three phases at \( T = 0 \). The error bars for \( T_{max} \), indicate 3 s.d. from fitting \( \rho(T) \) in Fig. 1a.

Figure 5 | Sketch of the interplay of vortex physics and quantum critical behaviour in the \( H–T \) phase diagram. Two critical fields, \( H_1^* \) and \( H_2^* \), are reported, separating three distinct phases at \( T = 0 \): a superconductor with \( \rho = 0 \) (dark blue) for all \( T < T_c(H) \) \( (T_c > 0) \) and \( H < H_1^* \), a superconductor with \( \rho = 0 \) only at \( T = 0 \) (that is, \( T_c = 0 \)) for \( H_1^* < H < H_2^* \), and an insulator (red), where \( \rho(T = 0) \to \infty \), for \( H_2^* < H \). The difference between the two superconducting ground states is in the ordering of the vortex matter: a pinned vortex solid (Bragg glass) for \( H < H_1^* \) and a vortex glass for \( H > H_1^* \), as shown schematically. The quantum critical regions (QCRs) corresponding to \( H_0^* \) and \( H_2^* \) are also shown schematically (dashed lines). The QC scaling associated with \( H_0^* \) does not extend to the lowest \( T \) (see dotted lines); apparently, this quantum phase transition is ‘hidden’ at low \( T \) by thermal fluctuations that cause the melting of the pinned vortex solid (Bragg glass) at \( T_m(H) \) for \( H < H_1^* \), and by the effects of disorder for \( H > H_1^* \). \( T_{max} \) is the temperature at which \( d\rho/dT \) changes sign.

At even higher \( H \), the transition from the superconducting vortex glass phase to an insulator with localized Cooper pairs occurs at \( H_1^* \) (where \( H_0^* < H_1^* < H_2^* \)), consistent with boson-dominated models of the SIT (ref. 13) and earlier arguments. The QC behaviour associated with this QPT is observed down to the lowest measured \( T \).

The scaling behaviour observed near critical fields \( H_1^* \) and \( H_2^* \) is consistent with a model where superconductivity is destroyed by quantum phase fluctuations in a 2D superconductor. We note that such scaling is independent of the nature of the insulator. In particular, the same scaling, except for the value of \( z \gamma \), is expected when the insulating phase is due to disorder and when it is caused by inhomogeneous charge ordering, which is known to be present in low-doped LSCO (refs 35–39), including charge-cluster glass and charge density waves.

We expect a similar two-stage \( H \)-field-tuned SIT to occur for other doping levels in the entire underdoped superconducting regime. This is supported by our observation of the same behaviour in samples that were prepared in two very different ways and in which the number of doped holes was probably not exactly the same.
Also, the possibility of two critical points was suggested recently for LSVO films with $x = 0.08, 0.09$ and $0.10$ (Supplementary Information). Our results provide important insight into the interplay of vortex line physics and quantum criticality. Apparently, the high-temperature, clean-limit SIT quantum critical fluid falls into the grip of disorder with decreasing $T$ and splits into two transitions: first, a $T = 0$ vortex lattice to vortex glass transition, followed by a genuine superconductor–insulator QPT at higher $H$.

**Methods**

**Samples.** The LSVO film with a nominal $x = 0.07$ was described in detail in ref. 11. The LSVO single crystal with a nominal $x = 0.06$ was grown by the travelling-solvent floating-zone technique\(^{22}\). Measurements were carried out on a sample that was cut out along the main crystallographic axes and polished into a sample with dimensions $3.02 \times 0.42 \times 0.34 \text{ mm}^3$ suitable for direct measurements of the in-plane resistance. Electrical contacts were made by evaporating gold on polished crystal surfaces, followed by annealing in air at 700 ºC. For current contacts, the entire area of two opposing side faces was covered with gold to ensure a uniform current flow through the sample. In turn, the voltage contacts were made narrow ($\sim 80 \mu\text{m}$) to minimize the uncertainty in the absolute values of the resistance. The distance between the voltage contacts is 1.41 mm. Gold leads were attached to the sample by using 6538 silver paste. This was followed by the heat treatment at 450 ºC in an oxygen flow for 6 min. The resulting contact resistances were less than 1 Ω at room temperature. As a result of annealing, the sample composition is probably La$_{1-x}$Sr$_x$CuO$_y$; with $y$ not precisely known.

**Measurements.** The in-plane sample resistance and magnetoresistance were measured with a standard four-probe a.c. method ($\sim 11$ Hz) in the Ohmic regime at current densities as low as $(3 \times 10^{-3} - 3 \times 10^{-4})$ A/cm$^2$ for the $x = 0.07$ film and $7 \times 10^{-4}$ A/cm$^2$ for the $x = 0.06$ crystal. To cover a wide range of $T$ and $H$, several different crystalotypes were used: a ‘He system at $T$ down to 0.3 K, with magnetic fields up to 9 T, using 0.02–0.05 T min$^{-1}$ sweep rates; a dilution refrigerator with $T$ down to 0.02 K and a ‘He system (0.3 K $\leq T \leq 30$ K) in superconducting magnets with fields up to 18 T at the National High Magnetic Field Laboratory (NHMFL), using 0.1–0.2 T min$^{-1}$ sweep rates; and a 35 T resistive magnet at the NHMFL with a variable-temperature insert (1.2 K down to 0.3 K), using 1 T min$^{-1}$ sweep rates. Therefore, the magnetoresistance measurements span more than two orders of magnitude in $T$, down to 0.09 K, which is much lower than the temperatures commonly employed in studies of underdoped cuprates. For that reason, we use excitations (that is, current densities) that are orders of magnitude lower than in similar studies. It was not possible to cool the samples below 0.09 K. The fields, applied perpendicular to the CuO$_2$ planes, were swept at constant temperatures. The sweep rates were always low enough to avoid any heating of the sample due to eddy currents. The resistance per square $R_{\text{CuO}_2} = \rho / d$ (d sample thickness; $\rho$, number of CuO$_2$ layers; $t = 6.6 \pm 0.8$ layer thickness of each layer); the resistance per square per CuO$_2$ layer $R_{\text{CuO}_2} = \rho / (d/\rho) = \rho / t$. 

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Author contributions

X.S. and D.P. conceived the project; the single crystal was grown by T.S.; X.S. and P.V.L. performed the measurements and analysed the data; V.D. and D.P. contributed to the data analysis and interpretation; X.S., P.V.L. and D.P. wrote the manuscript; D.P. planned and supervised the investigation. All authors commented on the manuscript.

Additional information

Supplementary information is available in the online version of the paper. Reprints and permissions information is available online at www.nature.com/reprints. Correspondence and requests for materials should be addressed to D.P.

Competing financial interests

The authors declare no competing financial interests.