**SUPERCONDUCTIVITY**

Restored strange metal phase through suppression of charge density waves in underdoped YBa$_2$Cu$_3$O$_{7-δ}$

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The normal state of optimally doped cuprates is dominated by the “strange metal” phase that shows a linear temperature ($T$) dependence of the resistivity persisting down to the lowest $T$. For underdoped cuprates, this behavior is lost below the pseudogap temperature $T^*$, where charge density waves (CDWs), together with other intertwined local orders, characterize the ground state. We found that the $T$-linear resistivity of highly strained, ultrathin, underdoped YBa$_2$Cu$_3$O$_{7-δ}$ films is restored when the CDW amplitude, detected by resonant inelastic x-ray scattering, is suppressed. This observation suggests an intimate connection between the onset of CDWs and the departure from $T$-linear resistivity in underdoped cuprates. Our results illustrate the potential of using strain control to manipulate the ground state of quantum materials.

Cuprate high-temperature superconductors belong to a class of materials where strong electron-electron correlations play a fundamental role, and whose unconventional properties might require abandoning traditional concepts of solid-state physics for a proper description (1). The “strange metal” phase of these superconductors is one of the most striking manifestations of the strong correlations. At optimal doping, this phase manifests as a linear temperature dependence of the resistivity that persists to the lowest $T$ when superconductivity is suppressed (1–4). This behavior is fundamentally different from that observed in more conventional metals (5), where a $T$-linear dependence of the resistivity is found only at high temperatures where phonon scattering dominates the transport. $T$-linear resistivity is also found in other strongly correlated systems, including pnictides (5) and magic-angle bilayer graphene (6).

Recent developments have attempted to model this behavior by considering that the scattering time approaches the fundamental Planckian limit defined by $\tau = \hbar/\hbar \nu T$ (where $\hbar$ and $\nu$ are the reduced Planck and Boltzmann constants) irrespective of the nature of the scattering process (4, 7). Local interactions among quasiparticles suggest that ordinary metals cannot thermalize on such a short time scale; reaching the Planckian limit would require every particle to be entangled with every other particle in the system.

In cuprates, the $T$-linear resistivity is lost for doping above and below the optimal doping. In the overdoped regime, the recovery of an almost $T^2$ dependence of the resistivity, typical of a Fermi liquid, is a consequence of the increased screening of the electron-electron interactions caused by a higher charge carrier density. In the underdoped regime, the deviation from $T$-linear behavior happens at temperatures close to $T^*$, known as the pseudogap temperature, where states are missing at the Fermi energy ($E_F$). In the pseudogap region, the high-temperature superconductor phase diagram also hosts a plethora of intertwined electronic local ordering phenomena that break rotational/translational symmetry (1, 9–13); charge density wave (CDW) order (14) is the most prominent one. The association between the departure from the $T$-linear resistivity and the occurrence of the pseudogap phenomenon has long been speculated. However, there is no consensus on the physics at play, nor on the causality hierarchy among pseudogap, local orders, and strange metal phenomenology (14, 15). This is because the strange metal exhibits its most salient features in transport and its connection to spectral signatures is elusive (16). More specifically, the challenge is to disentangle the various possible mechanisms leading to the breakdown of the $T$-linear resistivity, such as the occurrence of the pseudogap and the appearance of local orders such as CDW. One way to address this challenge is to tune the local properties of underdoped high-temperature superconductors. In particular, the CDW can be strongly modified under pressure (17, 18), strong magnetic fields (19), and strain in crystals (20) and thin films (21).

To tune the ground state in thin films of the cuprate material YBa$_2$Cu$_3$O$_{7-δ}$ (YBCO), we use the geometric modification of its unit cell under the strong strain induced by the substrate. We show that the $T$-linear resistivity dependence is completely recovered (down to the superconducting critical temperature $T_c$) along the $b$ axis when the CDW, detected by resonant inelastic x-ray scattering (RIXS), is suppressed along the $a$ axis.

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**Fig. 1.** Thickness dependence of the YBCO lattice parameters for films grown on MgO substrates. (A) Lattice parameters ($a$, $b$, and $c$ axes are indicated by squares, triangles, and circles, respectively) of YBCO, measured by x-ray diffraction at 300 K, as a function of the thickness of films with hole doping $p = 0.12$ grown on MgO. Thick lines are guides to the eye. (B) False-color scanning electron microscope image of a typical device used to measure the resistivity $\rho$, with the current $I$ flowing at an angle $\phi$ with respect to the YBCO [100] direction ($a$ axis). Inset: Overview of the patterned samples.

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Thin-film devices

The films we use span a wide range of hole-doping $p$, going from the strongly underdoped ($p = 0.10$) up to the optimally doped ($p = 0.17$) regime (22, 23). The strain was modified both by changing the substrate and by varying the film thickness $t$ in a range from 50 nm down to 10 nm. We used (100)-oriented MgO and 1° vicinal angle (001) SrTiO$_3$ (STO) substrates to grow untwinned YBCO films (24). When the YBCO thickness is reduced to a few unit cells ($t = 10$ nm), films grown on MgO are characterized by a considerable elongation of the $b$ axis and contraction of the $c$ axis, with the total volume of the unit cell unchanged with respect to relaxed systems (Fig. 1A). For films grown on vicinal cut STO, the YBCO unit cell is instead almost thickness-independent (fig. S1).

Figure 1B shows a typical device used to measure the resistivity as a function of temperature. The devices are patterned at an angle $\phi$ with respect to the YBCO [100] direction by using a carbon mask in combination with electron beam lithography and Ar$^+$ ion milling (22, 23, 25).

Angular dependence of in-plane resistivity

Figure 2A shows the temperature dependence of the resistivity $\rho$, measured in two devices oriented along the YBCO $a$ and $b$ axes ($\phi = 0^\circ$ and $90^\circ$, respectively) realized in an underdoped ($p = 0.11$) 50-nm-thick film grown on an MgO substrate. The resistivity anisotropy ratio at $T = 290$ K, defined by $\rho_b(290 \text{ K})/\rho_a(290 \text{ K})$, is ~1.2, a value fully compatible with the level of hole doping (26). The temperature $T_L$, estimated as the temperature below which the resistivity normalized to 290 K deviates by 1% from the linear fit ($\rho_a$ and $\gamma$ are the coefficients of the fit), $\rho_b(290 \text{ K})/\rho_a(290 \text{ K})$, is ~1.2. The dashed lines are the linear fits of the curves for $T > 260$ K. The inset shows the determination of $T_L$, which is the temperature where the resistivity normalized to 290 K deviates by 1% from the linear fit ($\rho_a$ and $\gamma$ are the coefficients of the fit). The temperature $T_L$ is determined as ~250 K.

The very different slopes of $\rho_b(T)$ and $\rho_a(T)$, patterned on a 10-nm-thick underdoped ($p = 0.11$) film. At $T = 290$ K, $\rho_b/\rho_a = 2.1$. For this kind of Fermi surface and the resistivity should be isotropic in the copper oxide planes. (E) Hypothetical anisotropic model Fermi surface that is compatible with the anisotropic transport of the 10-nm-thick devices.

(FIG. 2B, inset). Here, $\rho_a$ and $\gamma_a$ are respectively the intercept at $T = 0$ and the slope of the high-temperature linear dependence. At $p = 0.14$ and $p = 0.147$ (Fig. 2C and fig. S4C), for 10-nm-thick films we find that the $T$-linear behavior is completely recovered down to the superconducting transition. This finding is crucial because it indicates that the “strange metal” behavior is restored in ultrathin underdoped YBCO. What are the conditions for this to happen?

The most prominent structural effect we encounter by reducing the thickness of the films is that the YBCO $b$ axis expands, whereas the $a$ axis is only slightly modified; at the same time, the total volume of the cell remains unaltered as a consequence of a $c$-axis contraction. One of the effects of the strain is therefore to increase the orthorhombicity of few-nm-thick films. For 10-nm-thick underdoped films ($p = 0.12$), the values of the lattice parameters $a$ and $b$ are similar to those of YBCO with a much higher doping (29), close to the optimal doping. However, this effect cannot explain the anomalously high anisotropy ratio $\rho_b(290 \text{ K})/\rho_a(290 \text{ K})$ that we observe (Fig. 2B), a value in an optimally doped YBCO crystal is mainly related to the presence of CuO chains (26). Indeed, within a simple tight-binding description, an increase of ~0.02 Å in the $b$-axis lattice parameter—as we observe upon reducing the film thickness from 50 nm to 10 nm (Fig. 1A)—would reduce the corresponding hopping parameter between neighbor sites in the $b$ direction by only ~1% (30). Moreover, the resulting (weak) modification of the electronic structure would induce the opposite anisotropy in the $a$- and $b$-axis resistivities, namely $\rho_b > \rho_a$ (because of the larger $b$-axis value), relative to what we experimentally observe. This rules out the increased orthorhombicity as a possible direct explanation of the anisotropy we observe.

The very different slopes of $\rho_b(T)$ and $\rho_a(T)$ give a hint of the physics at play in the 10-nm-thick films. Following Boltzmann transport...
theory, the conductivity is given by

$$\sigma_{a,b}(T) = 2e^2 \sum \frac{v_{F,a,b}^2}{k} \frac{\Gamma(k)}{\Gamma(k')} (-n_{k'})$$  \hspace{1cm} (1)

where $v_{F,a,b}$ is the Fermi velocity, $\Gamma(k)$ is the $k$-dependent scattering rate, and $n_{k'}$ is the derivative of the Fermi distribution. Consequently, the ratio $\Gamma(k)/v_{F,a,b}$ determines the slope $\gamma_{a,b}$ of the temperature-dependent film resistivity. From the experimental data, we have $\gamma_a \gg \gamma_b$ (at $p = 0.11$, $\gamma_a = 4.9 \mu \Omega \cdot cm/K$ and $\gamma_b = 2.1 \mu \Omega \cdot cm/K$), which indicates that $v_{F,b} \gg v_{F,a}$ and therefore that the Fermi surface is already strongly anisotropic at room temperature. This is illustrated in Fig. 2, D and E, which show respectively the typical isotropic Fermi surface for the cuprates (where the contribution of the chains is neglected) and an anisotropic distorted Fermi surface, compatible with our experimental resistivity anisotropy. From Eq. 1, one may argue that anisotropic elastic scattering processes [e.g., due to small-angle scattering from impurities between CuO$_2$ planes (31)] can also account for the observed resistivity anisotropy. However, in this case $\Gamma(k)$ would be determined by the local density of states (32); that is, $\Gamma$ would still be proportional to $1/v_{F}$. Therefore, only an anisotropy of the Fermi velocity, which breaks the $C_4$ symmetry of the crystal, can account for the observed resistivity anisotropy.

A Fermi surface of the type shown in Fig. 2E can be the consequence of a strong electronic nematicity in the system. Such a state has been extensively investigated from a theoretical point of view (32–36) and it has been experimentally found, by in-plane resistivity anisotropy measurements, in tetragonal La$_2$-Sr$_2$CuO$_4$ films (37). We speculate that in the 10-nm thin films, the strain-induced distortion of the cell plays a fundamental role in stabilizing a nematic ground state already at room temperature.

The wider $T$-linear behavior of $\rho_b(T)$ in ultra-thin films is not observed in YBCO on vicinal angle STO substrates. Here $T_c$ does not change; it is the same along the $a$ and $b$ axes and going from 50-nm-thick to 10-nm-thick films (see fig. S3). This in agreement with the fact that the slopes $\gamma_{a,b}$ of $\rho_a(T)$ and $\rho_b(T)$ are comparable in 10- and 50-nm-thick films, so on STO substrates the Fermi surface is not substantially modified by strain at reduced thicknesses.

We have therefore arrived at the main result of our paper: A specific strain in underdoped 10-nm-thick YBCO films on a MgO substrate induces a nematic state that modifies the Fermi surface already at room temperature. But why would a distorted Fermi surface recover the $T$-linear resistivity behavior along the $b$ axis? The answer to this question comes from the study of the $\rho(T)$ dependence as a function of doping. Figure 3 shows $\rho_a(T)$ and $\rho_b(T)$ respectively as a function of the doping $p$ (Fig. 3, A and B) and of the angle $\phi$ (Fig. 3, C and D) for 10-nm-thick YBCO films on MgO substrates (as a function of the thickness in fig. S4).

The $\rho_b(T)$ curves are rather conventional and in agreement with previous results (22, 26). However, $\rho_p(T)$ (Fig. 3B) looks very different. As anticipated earlier, for $p = 0.14$ we have completely recovered the strange metal behavior: The $T$-linear dependence extends through the entire temperature range until superconductivity sets in. However, the situation changes at lower doping: For $p = 0.10$, the extracted $T_c$ along $b$ is close to the value along the $a$ axis (Fig. 3, C and D), and $\rho_b(T)$ shows a pronounced upturn at low temperature before the superconducting transition. The upturn of the $\rho_b(T)$ in our 10-nm-thick YBCO films is observed at higher doping relative to the $\rho_a(T)$ (Fig. 3, A and B); indeed, it already appears at $p = 0.13$. This upturn in the resistivity has been attributed to the loss of high-mobility electron pockets because of the CDW order ending at $p = 0.08$ (38) and/or to the proximity to the antiferromagnetic Mott insulator through spin density waves (39–41). The doping dependence of $\rho_b(T)$ therefore points toward a strong involvement of the CDW order in the phenomenon we observe.

Resonant inelastic x-ray measurements of CDW

To characterize the CDW order in the YBCO films, we used RIXS at the Cu $L_3$ edge (~930 eV) (23). We investigated films with thickness of 10 and 50 nm grown on both MgO and STO at the doping $p = 0.125$ where the intensity of the CDW is strongest (17, 42). To isolate the contribution of the CDW, we measured RIXS spectra at $T = 70$ K (i.e., close to $T_c$), where the CDW signal is maximized, and at $T = 200$ K, a temperature where the CDW contribution is negligible. The CDW peak has been explored along $a$ and $b$ with two orthogonal cuts along the $H$ and $K$ directions of reciprocal space, centered around the wave vector of the CDW $q_{\text{CDW}} = (H_{\text{CDW}}, K_{\text{CDW}})$. Shown in Fig. 4, A and B, is the quasi-elastic component of the RIXS spectra along the $a$ and $b$ axes for a 50-nm-thick film on MgO, as a function of $H$, at $T = 70$ K and $T = 200$ K. Along both directions, at $T = 200$ K only a broad peak is present (Fig. 4, A

Fig. 3. Angular dependence of the in-plane resistivity of the 10-nm YBCO on MgO as a function of oxygen doping. (A) $\rho(T)$ measured along the a-axis direction of 10-nm-thick films, with different doping levels, grown on MgO substrates. $T_c$ is extracted as described in Fig. 1. The hatched regions represent the temperature range above $T_c$ where at each $p$ the $T$-linear resistivity regime occurs. (B) Same as (A), but with $\rho(T)$ measured along the b-axis. (C) $\rho(T)$, normalized to its value at $T = 290$ K, for a 10-nm-thick film ($p = 0.120$) on MgO is shown as a function of $\phi$ (i.e., with respect to the a-axis direction). For each angle, the temperature $T_c$ has been extracted (black spheres and dots). $T_c$ is approximately constant as a function of $\phi$, except around $\phi = 90^\circ$, where it exhibits a suppression. (D) $T_c$ values as a function of $\phi$ for 10-nm-thick films with different doping levels ($p = 0.099, 0.103, 0.108, 0.117, 0.118, 0.120, 0.123, 0.134, 0.140, or 0.147$).

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and B, red regions); at $T = 70$ K, the signal is given by the sum of a broad peak (similar to that measured at high temperature) and of a narrow peak. This narrow, temperature-dependent peak is a signature of CDW; the broad, almost temperature-independent peak is instead a signature of short-range charge density fluctuations (CDFs) (43) (i.e., charge modulations, precursors of CDWs), which are present in the phase diagram at any temperature and in a very broad doping range, including the overdoped region. The same measurements as in Fig. 4, A and B, are reported in Fig. 4, C and D, for a 10-nm-thick film on MgO. Here, along the $a$ axis (Fig. 4D), both the CDW and the CDF peaks are very similar to those measured in the thick film. Along the $a$ axis (Fig. 4C), the situation is markedly different. The broad-in-$q_y$/CDF peak is unchanged with respect to the 50-nm-thick sample; in contrast, the narrow CDW peak, emerging at lower temperature, is almost negligible. This occurrence has been verified on the same sample, measuring along the $K$ direction (fig. S5), and on other ultrathin films on MgO with different doping (fig. S6).

Note that our films are not perfectly untwinned (the untwining degree is ~85% for films grown on MgO, ~90% for films grown on STO). The small temperature-dependent CDW signal measured along the $a$ axis in the 10-nm films on MgO can be attributed to twinned domains. Indeed, once the twinning is taken into account, the actual CDW signal becomes effectively negligible in the $a$-axis direction. We conclude that in 10-nm-thick films on MgO, the CDW is unidirectional and directed along the $b$ axis, whereas the CDW in YBCO grown on STO is thickness-independent (fig. S7). Strain-induced modifications of CDW have also been recently observed in YBCO single crystals, where the in-plane uniaxial compression along either $a$ or $b$ gives rise to an enhancement of the CDW in the orthogonal direction (44).

**Discussion**

Our results are in line with the theoretical predictions of unidirectional CDWs in the presence of nematicity (34, 35), as a consequence of the modified Fermi surface. Our experiment additionally shows a clear correlation between the $T$-linear dependence of $p(T)$ and the charge order: The recovery of the $T$-linear resistivity along the $b$ axis is associated with the suppression of the CDW along the YBCO $a$ axis. In this scenario, only CDFs survive in the system at any temperature, which might be relevant both for Planckian metal (45) and marginal Fermi liquid theories (46) of the strange metal.

The strong correlation between the CDW and the breakdown of the $T$-linear behavior is further supported by the coincidence of the onset temperatures for the two phenomena. Figure S8 shows the temperature dependence of the CDW order for our 50- and 10-nm-thick films on MgO and STO substrates. Within experimental error, $T_{\text{CDW}}$ (for $p = 0.125$) is the same as the temperature at which the resistivity departs from the linearity $T^*$ and the same as the pseudogap temperature $T^*$ [taken from literature (8, 27)]. This is illustrated in Fig. 5, which shows a revised phase diagram for the 10 nm-thick films. Our RIXS measurements therefore require the revision of the common belief that the CDW is a low-temperature phenomenon happening at temperatures well below $T_L$ and $T^*$, as indicated in early studies (II, 42). The coincidence among $T^*$, $T_L$, and $T_{\text{CDW}}$ is a common feature of doping above $p = 1/8$ where the CDW order is strong, as also supported by recent RIXS data (43), which are included for completeness in Fig. 5.

The quality of our samples allows us to exclude the possibility that structural changes of the CuO chains along the $b$ axis in 10-nm films might have a role in the phenomenology we have observed. For overdoped films, where the conductivity of the chains causes an upward deviation from linearity in $\rho(T)$, we observe the same $\rho(T)$ dependence for 50-nm and 10-nm films and very similar values of $T_L$ (fig. S2), which excludes CuO chain modification effects in our very thin films.

However, there is still an issue to be addressed: Why, in our films, does the disappearance of the CDW along the YBCO $a$ axis lead to a recovery of the $T$-linear dependence of the resistivity along the $b$ axis (Figs. 2B and 4A)?

**Fig. 4. Thickness dependence of the CDW in YBCO thin films.** (A and B) Quasi-elastic scans measured at $T = 70$ K and $T = 200$ K on the 50-nm thin film on MgO along the $(H, 0)$ direction (i.e., $a$ axis) (A) and the $(H, k_{\text{CDW}})$ direction (i.e., $b$ axis) (B). (C and D) Same as (A) and (B), but on a 10-nm-thick sample. The measurements were performed along the $(H, 0)$ direction (i.e., $a$ axis) (C) and the $(H, k_{\text{CDW}})$ direction (i.e., $b$ axis) (D). In the 10-nm-thick sample, the CDW intensity along the $a$ axis is almost negligible. If we take into account the percentage of twin domains (~15%) present in our films, we conclude that in films a few unit cells thick on MgO, the CDW is unidirectional along the $b$ axis. The green line in the inset of each panel shows the direction of the scan relative to the CDW peak in reciprocal space.
strongly affects the resistivity along the a axis while leaving the resistivity along the b axis unaltered.

Finally, we consider the role of the pseudogap in the departure of the CDW-linear resistivity in underdoped cuprates. From the slopes of $\rho(T)$ along the $a$ and $b$ axes, $\rho_{a,b}$, we infer that the Fermi surface is anisotropic at any doping (the ratio $\rho_{a}/\rho_{b}$ is almost doping-independent). For $p \leq 0.1$, we observe that $T_c$, which at that doping level is much higher than $T_{CDW}$ and comparable to $T^*$ (taken from literature ($8, 27$)), is isotropic ($T^* = T_c^*$). This hints at a pseudogap that is isotropic on a distorted Fermi surface. Such an isotropic pseudogap cannot explain the otherwise anisotropic $T_c$ we observed at higher doping ($p > 0.1$). This allows us to conclude that the pseudogap is not the main effect responsible for the departure of the linearity when the CDW order is strong ($p > 0.11$), which is the range of doping in our experiment. However, one cannot exclude a major influence of the pseudogap on the transport at lower doping when $T^*$ becomes much higher (close to room temperature and beyond) than $T_{CDW}$.

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**Supplementary Materials**

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Materials and Methods

Supplementary Text

Figs. S1 to S9

References (49–64)

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