Superfluid stiffness of a KTaO₃-based two-dimensional electron gas

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After almost twenty years of intense work on the celebrated LaAlO₃/SrTiO₃ system, the recent discovery of a superconducting two-dimensional electron gas (2-DEG) in (111)-oriented KTaO₃-based heterostructures injects new momentum to the field of oxides interface. However, while both interfaces share common properties, experiments also suggest important differences between the two systems. Here, we report gate tunable superconductivity in 2-DEGs generated at the surface of a (111)-oriented KTaO₃ crystal by the simple sputtering of a thin Al layer. We extract the superfluid stiffness of the 2-DEGs and show that its temperature dependence is consistent with a node-less superconducting order parameter having a gap value larger than expected within a simple BCS weak-coupling limit model. The superconducting transition follows the Berezinskii-Kosterlitz-Thouless scenario, which was not reported on SrTiO₃-based interfaces. Our finding offers innovative perspectives for fundamental science but also for device applications in a variety of fields such as spin-orbitronics and topological electronics.

Potassium tantalate KTaO₃ is a band insulator with a 3.6 eV gap that retains a cubic perovskite structure down to the lowest temperature¹. Like strontium titanate (SrTiO₃), it is a quantum paraelectric material on the verge of a ferroelectric instability that is characterized by a large permittivity at low temperature (εₑ ≃ 5000)². Both materials can be turned into metal by electron doping, through oxygen vacancies, for example. Because of their common properties, it was suggested that superconductivity should also occur in doped KTaO₃. However, while superconductivity was discovered more than half a century ago in bulk SrTiO₃³, all the attempts to induce bulk superconductivity in KTaO₃ have failed so far⁴. Using ionic gating, Ueno et al. could generate a superconducting 2-DEG at the surface of (001)-KTaO₃, albeit at a very low temperature (Tₑ ≃ 40 mK)⁵. Later explorations of KTaO₃ 2-DEGs did not evidence any superconductivity until the beginning of the year 2021, when two articles reported the discovery of superconducting 2-DEG formed at the interface between (111)-KTaO₃ and insulating overlayers of LaAlO₃ or EuO⁶,⁷. An empiric increase of Tₑ with electron density was proposed with a maximum value of 2.2 K for doping of ≃ 1.04 × 10¹⁴ e⁻ cm⁻²⁶, which is almost one order of magnitude higher than in the LaAlO₃/SrTiO₃ interface⁸. An electric field effect control of the Tₑ was also demonstrated in a Hall bar device⁷ and a dome-shaped superconducting phase diagram similar to that of SrTiO₃-based interfaces was derived⁹,¹⁰. Following this discovery, the (110)-oriented KTaO₃ interface was also found to be superconducting with Tₑ ≃ 1 K¹¹. It was recently proposed that the soft transverse optical mode, involved in the quantum paraelectricity, could be responsible for electron pairing in KTaO₃ interfaces. The coupling amplitude between this phonon mode and electrons is expected to be maximum in the (111) orientation and minimum in the (001) one, which would explain the hierarchy in Tₑ observed in these superconducting 2-DEGs¹².

In conventional superconductors, well described by the Bardeen–Cooper–Schrieffer (BCS) theory, the superconducting transition is controlled by the breaking of Cooper pairs as the temperature exceeds the energy scale set by the superconducting gap.

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However, in two-dimensional superconductors, the superfluid stiffness, i.e., the energy associated with the phase rigidity of the superconducting condensate, can be comparable to the pairing energy, allowing for a $T_c$ suppression driven by the loss of phase coherence. In this case, the transition is expected to belong to the Berezinskii–Kosterlitz–Thouless (BKT) universality class, where the transition is controlled by the unbinding of topological vortex–antivortex pairs\(^\text{13–15}\). Critical magnetic field measurements in (111)-KTo$O_3$ 2-DEGs, both in the perpendicular and in the parallel geometry, set an upper bound, $d = 5$ nm, on the extension of the 2-DEG in the substrate\(^\circ\). This is lower than the superconducting coherence length, $\xi = 10–15$ nm\(^\circ\), which confirms that the superconducting 2-DEG is within the 2D limit. In addition, the presence of disorder, which has been identified in this system\(^\circ\), is also expected to lower the superfluid rigidity and reinforce the role of phase fluctuations. Even though the measurements of the current-voltage characteristics in ref. 6 could be compatible with indirect signatures of a BKT transition, a direct measurement of the superfluid stiffness is required to properly address this issue\(^\circ\).

Here, we show that a 2-DEG can be generated at the surface of a (111)-oriented KTo$O_3$ crystal simply by sputtering a very thin Al layer. The deposition of Al leads to the reduction of Ta ions as evidenced by X-ray photoelectron spectroscopy (XPS) and leads to the formation of an interfacial gate-tunable superconducting 2-DEG. We use resonant microwave transport to measure the complex conductivity of the 2-DEG and extract the temperature-dependent superfluid stiffness $f_s(T)$. Our results are consistent with a node-less superconducting order parameter in a rather strong-coupling regime $(\Gamma(0)/k_B T_c = 2.3)$. Taking into account the presence of disorder and finite-frequency effect, we show that the superconducting transition follows the Berezinskii–Kosterlitz–Thouless model, which was not observed on SrTo$O_3$-based interfaces.

Results

2-DEGs were generated by dc sputtering of a very thin Al layer on (111)-oriented KTo$O_3$ substrates at a temperature between 550 and 600 °C. The preparation process is detailed in the Methods section. Prior to deposition, we measured the in-situ X-ray photoelectron spectra (XPS) of the Ta 4f valence state (Fig. 1a top) of the KTo$O_3$ substrate. The spectra show the sole presence of Ta$^{3+}$ states (4f$^{12}$ and 4f$^{12}$), indicating the expected stoichiometry of the substrate. The Ta 4f core levels were then measured after growing 1.8–19 nm of Al and transferring the sample in vacuum to the XPS setup. The bottom graph in Fig. 1a shows the Ta 4f core level spectra with additional peaks corresponding to reduced states of Ta, i.e., Ta$^{2+}$ and Ta$^{4+}$. The deeper and lighter shades of the same colors represent the 4f$^{12}$ and 4f$^{12}$ valence states of the respective peaks. The data and sum fit envelope are shown in red circles and black lines. EELS maps (Al-L$\text{_{2,3}}$, Ta-O$\text{_{2,3}}$, and K-L$\text{_{2,3}}$, edges) showing the presence of Al on top of the interface without any interdiffusion but with some limited diffusion of Ta and K inside the AlO$X$ layer.

![image](https://example.com/image.png)

**Fig. 1** XPS and STEM characterization of the AlO$X$/KTo$O_3$ samples. a X-ray photoelectron spectra near the 4f core level binding energy of Ta for a KTo$O_3$ substrate prior to deposition (top) and after deposition of 1.8 nm of Al (bottom). The fitted peaks for Ta$^{3+}$, Ta$^{2+}$, and Ta$^{4+}$ are shown in cyan, green and violet colors, respectively. The deeper and lighter shades of the same colors represent the 4f$^{12}$ and 4f$^{12}$ valence states of the respective peaks. b (top) HAADF scanning transmission electron microscopy image at the cross-section of KTo$O_3$ (111) and AlO$X$ interface. The KTo$O_3$ is observed along the [112] direction. (down) EELS maps (Al-L$\text{_{2,3}}$, Ta-O$\text{_{2,3}}$, and K-L$\text{_{2,3}}$, edges) showing the presence of Al on top of the interface without any interdiffusion but with some limited diffusion of Ta and K inside the AlO$X$ layer.

In the present work, four samples, labeled A, B, C, and D, have been investigated by transport measurement at low temperature in a dilution refrigerator (see Methods section for fabrication parameters).

Figure 2a shows the resistance vs temperature curve of sample A on a wide temperature range revealing a superconducting transition at $T_c = 0.9$ K. In Fig. 2b, we plot the $T_c$ as a function of the 2D carrier density, $n_{2D}$ for the different samples studied and compare their values with those extracted from ref. 6. Our results confirm the trend observed in the literature ($T_c$ increases with the carrier density) and demonstrate that our growth method, while being much easier to implement than the molecular beam epitaxy of a rare-earth element such as Eu or the pulsed laser deposition of a complex oxide, is able to produce good quality samples with similar $T_c$. The resistance vs temperature curves of sample B measured for different values of a magnetic field applied perpendicularly to the sample plane are shown in Fig. 2c. The temperature dependence of the critical magnetic field is consistent with a Landau-Ginsburg model near $T_c$, $\mu_B H_c(T) = \mu_B H_c(0) (1 - T/T_c)^3$. We found $H_c(0) = 27$ nm, which is comparable with the value reported in ref. 6.

Although KTo$O_3$ is a quantum paraelectric material like SrTo$O_3$, its permittivity is reduced by a factor of five as compared to SrTo$O_3$, making the electric field effect less efficient in a back-gating configuration\(^2\). To overcome this difficulty, we prepared an AlO$X$/KTo$O_3$ sample using a thinner substrate (150 µm). After cooling the sample, the back-gate voltage was first swept to its maximum value $V_g = 200$ V while keeping the 2-DEG at the electrical ground. This forming procedure is commonly applied on SrTo$O_3$-based interfaces to ensure the reversibility of the gate sweeps in further gating sequences\(^3\). Figure 3 shows the sheet resistance of sample C as a
function of temperature for different values of the gate voltage between $-40$ and $200$ V. Electrostatic gating induces both a modulation of the normal-state resistance and a variation of the superconducting critical temperature. For negative gate voltages corresponding to a depleted quantum well, $R$ vs $T$ curves exhibit a quasi-reentrant behavior: the resistance first decreases and then upturns upon further cooling$^{22,23}$. This is characteristic of disordered superconducting thin films in which superconductivity only exists locally, forming a network of isolated islands surrounded by an insulating medium that precludes percolation. While the decrease of resistance marks the emergence of superconductivity inside the islands, the upturn of resistance at low-temperature results from the opening of a gap in the excitation spectrum, which prevents the flow of quasiparticles across islands. Hence, the resistance does not reach zero, indicating that the superconducting order does not extend at a long range. As carriers are added upon increasing the gate voltage, the resistance curves flatten at low temperature, and the 2-DEG eventually reaches a true zero resistive state ($V_G > -25$ V).

**Fig. 2 | Magneto-transport characterization.** a Sheet resistance of sample A as a function of temperature (log scale) showing a superconducting transition at $T_c \approx 0.9$ K. b Summary of the superconducting critical temperature as a function of the carrier density for all samples studied in this work compared with results from ref. 6. c Sheet resistivity of sample B as a function of temperature for increasing magnetic fields between 0 and 750 mT. d Perpendicular critical magnetic field defined by a 50% drop of the normal-state resistance measured at $T = 1.65$ K. The orange line corresponds to a linear fit with the Ginsburg Landau formula.

**Fig. 3 | Electric field effect and superconducting phase diagram.** a Temperature dependence of the sheet resistance of sample C for different values of the back-gate voltage in the range $-40$ to $200$ V. b Normalized sheet resistance in color scale as a function of the carrier density extracted from Hall measurements and temperature. The dashed line indicates the critical temperature defined by an 80% drop of the resistance with respect to the normal resistance $R_N$ at $T = 0.9$ K. c Sheet resistance at $0.9$ K as a function of the back-gate voltage.
A superconducting order is established through Josephson coupling between the islands. Further doping makes the network of islands denser and increases the coupling between islands resulting in a “homogeneous-like” superconducting 2-DEG at high doping. The resulting superconducting phase diagram is shown in Fig. 3b, where the resistance is plotted in color scale as a function of temperature and electron density extracted by combining the Hall effect and gate voltage by more than one order of magnitude. This indicates that the gate voltage not only controls the carrier density but also modifies deeply the electronic properties of the 2-DEG, in particular the electronic mobility, in agreement with the previous reports..

We further investigated the superconducting KTaO3 2-DEG by measuring its superfluid stiffness $\kappa$, which is the energy scale associated with the phase rigidity of the superconducting condensate. $\kappa$ is related to the imaginary part of the complex conductivity $\sigma(\omega) = \sigma_i(\omega) - i\sigma_f(\omega)$ of the superconductor that accounts for the transport of Cooper pairs at finite-frequency $\omega$. This is a direct probe of the superconducting order parameter that provides important information on the nature of the superconducting state. In the low-frequency limit, $\hbar\omega \ll \Delta$, a superconductor displays an inductive response to an ac electrical current and $\sigma_f(\omega) = \frac{\hbar}{2eL_s(I)}$, where $L_s$ is the kinetic inductance of the superconductor that diverges at $T_c$. The superfluid stiffness is then directly related to $L_s$

$$\kappa(T) = \frac{\hbar^2}{4e^2L_s(T)} \tag{1}$$

where $\hbar$ is the reduced Planck constant and $e$ is the electron charge.

We used resonant microwave transport to extract $L_s$ below $T_c$ and determine the superfluid stiffness of sample D as a function of temperature. The method, which was successfully applied to superconducting SrTiO$_3$-based interfaces, is illustrated in Fig. 4a and described in detail in refs. 26,27. In short, the KTaO$_3$ sample is embedded into a parallel RLC resonant electrical circuit made with surface mount microwave devices (SMD). The capacitance of the circuit is dominated by the KTaO$_3$ substrate contribution (C$_{KTO}$) due to its large intrinsic permittivity. The total inductance of the circuit, $L_{tot}(T) = L_{s(M)}(T) + L_{KTO}$, includes the contribution of an SMD inductor ($L_M$) and the contribution of the kinetic inductance $L_s$ of the superconducting 2-DEG below $T_c$. Finally, an SMD resistor $R_1$ imposes that the dissipative part of the circuit impedance remains close to 50 $\Omega$ in the entire temperature range, ensuring a good impedance matching with the microwave circuitry.

The circuit resonates at the frequency $\omega_0 = \frac{1}{\sqrt{L_{tot}(T)C_0}}$, which is accessed by measuring the reflection coefficient of the sample circuit $\Gamma(\omega) = \frac{\rho_{in}}{\rho_{out}} = \frac{Z(\omega)}{-Z(\omega)}$, where $Z(\omega)$ is the impedance of the sample for any frequency and $\rho_{in}$, $\rho_{out}$ are the incident and outgoing waves.

The resonance manifests itself as a peak in the real part of the circuit impedance, $Z(\omega)$, accompanied by a $\pi$ phase shift. The height and the width of the peak are controlled by the dissipative part of the circuit impedance. In the superconducting state, the 2-DEG conductance acquires a kinetic inductance $L_s$ that generates a shift of $\omega_0$ towards high frequencies with respect to the normal state (Fig. 4b). The temperature-dependent superfluid stiffness $\kappa^{exp}$, extracted from the resonance shift and Eq. (1) is presented in Fig. 4c (blue circles).

### Discussion

The flattening of the $\kappa^{exp}$ curve below 1 K supports a fully gapped behavior, i.e., an absence of nodes in the order parameter. The purpled dashed line ($\kappa^{BCS}$) shows an attempt to fit the experimental curve with a standard BCS expression $\kappa^{BCS}(T) = (\Delta(T)/\Delta(0))\tanh(\Delta(T)/k_B T_c)$, where $\Delta(T)$ is the superconducting gap obtained numerically by a self-consistent solution of the BCS equation, so that it vanishes at the mean-field temperature $T^0_c$ (i.e., the temperature at which Cooper pairs form). Since $\Delta(0)$ is fixed by the experimental value at the lowest temperature, the only free parameter is then the ratio $\Delta(T=0)/k_B T_c$, that determines the curvature of the $\kappa(T)$ curve. As one can see, even using a relatively strong coupling value $\Delta(0)/k_B T^0_c = 2.3$, from the fit of the low-temperature curve, one obtains $T^0_c = 2.2$ K, that is larger than the experimental $T_c$. To fit the data in the whole temperature range with the BCS expression only, one would then need an unreasonably large value $(\Delta(0)/k_B T^0_c = 6)$, a result that we checked to hold irrespectively of the exact functional BCS form used to fit the stiffness.
We here follow a different approach and interpret the rapid drop of \( J_0(T) \) below the BCS fit as a BKT signature, as we will discuss below. This interpretation is supported by a second striking observation that holds regardless of any specific consideration about its temperature dependence: the \( T = 0 \) value of the stiffness \( J_0(T=0) \approx 7.3 \, \text{K} \) is of the same order as \( T_c \approx 2.2 \, \text{K} \). It is worth noting that in conventional superconductors, where the superfluid density \( n_s(T=0) \) is close to the carrier density \( n_{2D} \), the stiffness at zero temperature is of the order of the Fermi energy, and then several orders of magnitude larger than \( T_c \). A strong reduction of \( J_0(T) \) is instead observed in 2D superconductors, where disorder strongly reduces \( n_s \) with respect to \( n_{2D} \) already at \( T = 0 \). In the dirty limit, in which the elastic scattering rate \( 1/\tau \) is much larger than the superconducting gap, only a fraction of carriers, \( n_s/n_{2D} \geq \Delta(0)/(k_BT_c) \), forms the superconducting condensate. In a single-band picture, an estimate of the superfluid stiffness is obtained from \( \Delta(0) \) and the normal resistance \( R_N \), \( J_0 \approx \frac{k_BT_c}{2\pi \hbar} \). Using the previously estimated value of \( \Delta(0) = 5 \, \text{K} \) and \( R_N = 1300 \, \Omega \), we obtain \( J_0 \approx 11.8 \, \text{K} \), close to the measured value \( (f^{\text{PD}}(T=0) \approx 7.3 \, \text{K}) \), which is consistent with the dirty limit.

The superfluid density of the 2-DEG can be directly deduced from the stiffness through the formula \( n_s = \frac{4m}{\pi} J_0 \), where \( m \) is the effective mass of superconducting electrons. In the case of (111)-KTaO3 2-DEGs, the condensation field is derived from the bulk 3-2 states with a Fermi surface formed by a hexagonal contour inside a sixfold symmetric star-shaped contour. Considering an average effective mass \( 0.5m_0 \), the corresponding superfluid density \( n_s \) extracted from \( f^{\text{PD}} \) is \( n_s = 1.8 \times 10^{12} \, \text{cm}^{-2} \), which is about 2.5% of the total carrier density \( (n_{2D} = 7.5 \times 10^{13} \, \text{cm}^{-2} \) for sample D). This very low ratio is comparable with previous findings in LaAlO3/SrTiO3 interfaces. Although such reduced superfluid density is consistent with the dirty limit, KTaO3 (111) 2-DEG is a multiband system, in which superconductivity may involve only specific bands, as also suggested in SrTiO3.

The reduced dimensionality and the suppression of the energy scale associated with the stiffness represent the prerequisites to observe BCS-like physics, since it makes the BKT temperature scale \( T_{\text{BKT}} \) associated with the unbinding of vortex-antivortex pairs far enough from \( T_c \). The most famous hallmark of the BKT transitions is the discontinuous jump to zero of \( J_0 \) at \( T_{\text{BKT}} \approx T_c \), with a universal ratio \( J_0(T_{\text{BKT}})/J_0(T_c) = 2/\pi^4 \). Such a prediction, theoretically based on the study of the 2D XY model, has been successfully confirmed in superfluid He films. In practice, the experimental observation of the BKT transition in real superconductors is more subtle. Indeed, in thin films, the suppression of \( n_s \), and then \( J_0 \) with disorder comes along with an increasing inhomogeneity of the SC background, that is predicted to smear out the discontinuous superfluid-density jump into a rapid downturn, as observed experimentally via the direct measurement of the inverse penetration depth or indirectly via the measurement of the exponent of the nonlinear IV characteristics near \( T_c \). In the case of SrTiO3-based interfaces, the direct measurement of \( J_0 \) is rather challenging, and the few experimental reports available so far do not evidence a BKT jump.

Within the BKT approach, the effect of vortex-like topological excitations provides additional suppression of \( J_0 \), with respect to the BCS dependence discussed above, driven only by quasiparticle excitations. To provide a fit of \( f^{\text{PD}} \), we then solved numerically the renormalization-group (RG) equations of the BKT theory for the superfluid stiffness and vortex fugacity. As input parameters of the RG equations, we used the BCS temperature dependence of the stiffness. As mentioned above, the low-temperature part is fully captured by the BCS approximation, and for the estimated \( \Delta(0)/k_BT_c \) ratio, the dirty-limit and the clean-limit expressions of \( f^{\text{BCS}} \) provide the same result. We also included the finite-frequency effects in our calculation. Indeed, even though the resonance frequency (about 0.5 GHz) is still small as compared to the optical gap (2A - 10 K - 200 GHz) it can nonetheless lead to non-negligible effects, in particular a rounding of the jump and suppression of the stiffness at a temperature slightly larger than the one where the dc resistivity vanishes as observed in our case. The resistivity itself is consistently fitted with the interpolating Halperin–Nelson formula, which accounts for BKT-like fluctuations between \( T_{\text{BKT}} \) and \( T_c \), and for standard Gaussian fluctuations above \( T_c \). Finally, to account for spatial inhomogeneities, we introduce a gaussian distribution of local \( T_c \) and \( J_0 \), with variance \( \sigma_T \) centered around \( T_{\text{BKT}} \) and \( f^{\text{PD}}(0) \). As seen in Fig. 4c, the result of the fitting procedure (dashed red line) is in very good agreement with experimental data both for the superfluid stiffness and the resistance, considering a very small inhomogeneity, \( \sigma_T = 0.02 \). Details on the fitting procedure are given in the Method Section.

Although KTaO3 and SrTiO3 have many common properties, the superconducting phases of their interface 2-DEG exhibit noticeable differences. Whereas a pure BCS weak-coupling limit with \( \Delta(0)/k_BT_c = 1.76 \) provides a very good description of superconductivity in SrTiO3-based interfaces, we found a stronger value of the coupling for KTaO3 (\( \Delta(0)/k_BT_c = 2.3 \)). Such an important difference, which must be traced back to the pairing mechanism, is a strong constraint on the possible origin of superconductivity in these two materials. In addition, BKT physics was not observable in SrTiO3 for which a simple BCS model without phase fluctuations was sufficient to fit the \( J_0(T) \) curves with a very good accuracy. This may suggest more bosonic-like superconductivity in KTaO3-based interfaces (in the highly doped regime), as evidenced by the large separation between the pairing scale, set by \( \Delta \), and the phase-coherence scale, set by the small value of the superfluid stiffness. Recent measurements of the in-plane critical field in KTaO3-based interfaces suggested that the order parameter could be a mixture of s-wave and p-wave pairing components induced by strong spin-orbit coupling. While we can not rule out this possibility, the saturation of the \( J_0(T) \) curve below \( T_{\text{BKT}} \) seen in Fig. 4b suggests a dominance of the fully gapped s-wave component. In addition, despite the complex band structure of the KTaO3(111) interfaces, we have not observed any signatures suggesting multi-gap superconductivity in our data. Further experiments, including tunneling spectroscopy, are therefore necessary to understand the nature of superconductivity in KTaO3-based interfaces.

### Methods

#### Sample fabrication

Prior to deposition, KTaO3 (111) substrates from MTI corporation were annealed at 600 °C for 1h in vacuum. Then, the thin Al layer was deposited in a dc magnetron sputtering system (PLASSYS MP450S) under a base pressure of the vacuum chamber lower than 5 × 10⁻⁶ mbar. During Al deposition, the Ar partial pressure and the dc power were kept fixed at 5 × 10⁻⁸ mbar and 10 W, respectively. The deposition rate for Al was 0.66 A/s. Table 1 below summarizes the deposition parameters for the different samples.

#### XPS analysis

X-ray photoelectron spectroscopy (XPS) was measured using a non-monochromatized Mg Kα source (hv = 1253.6 eV) in an Omicron NanoTechnology GmbH system with a base pressure of 5 × 10⁻¹⁰ mbar. The operating current and voltage of the source was 20 mA and 15 kV, respectively. Spectral analysis to determine different valence states of Ta were carried out using the CasaXPS software. Adventitious carbon was used as a charge reference to obtain the Ta 4f₅/₂ peak position for the fitting. The energy difference and the ratio of the area between 4f₅/₂ and 4f₇/₂ peaks for all the Ta valence states were constrained according to the previously reported values.

#### STEM characterization

STEM-HAADF and STEM-EELS measurements have been done at 100 kV using a Cs corrected Nikon STEM microscope and a Gatan modified EELS spectrometer equipped with a MerlinEM detector.
Theoretical analysis of $J_f(T)$

In order to account for vortex excitations, we solved the BKT RG equations \(^{15,43}\) for the vortex fugacity $g = 2\pi e^{-\mu/\sqrt{4\pi}}$, with $\mu$ the vortex-core energy, and the rescaled stiffness $k = n^2/|k_B T|$:

$$\frac{dK}{dt} = -K^2 g^2, \quad (2)$$

$$\frac{dg}{dt} = (2 - K) g. \quad (3)$$

where $t = \ln(t/\xi_0)$ is the RG-scaled lattice spacing with respect to the coherence length $\xi_0$ that controls the vortex sizes and appears as a short-scale cut-off for the theory. The initial values at $t = 0$ are set by the BCS fitting $\mu_{BCS}(T)$ of $\mu_{BCS}$, and the renormalized stiffness is given by the large-scale behavior, $J_f = (k_B T/\ln (\xi/\xi_0))$. The ratio $\mu_{BCS}/T = 0.87$, similar to the one found in other conventional superconductors \(^{16,35,37,38}\), is used as a free (temperature-independent) parameter, which controls the strength of stiffness renormalization due to bound vortices below $T_{BKT}$ \(^{39}\). To account for finite-frequency effects, we further include a dynamical screening of vortices \(^{49,50}\) via an effective frequency-dependence dielectric function $\epsilon(\omega)$ which enters in the complex conductivity of the film as $\sigma(\omega) = \frac{i}{4\pi} \omega \epsilon(\omega)$. At zero frequency $\epsilon(\omega)$ is real and $\epsilon(0) = 1$, while $\epsilon(\omega)$ is complex. To recover the usual static result. At finite-frequency $\epsilon(\omega)$ develops an imaginary part due to the vortex motion, that can be expressed in the first approximation \(^{39}\) as $\epsilon_2 = (r_2/\xi_0^2)$, where $\xi$ is the vortex correlation length and $r_2$ is a finite length scale set in by the finite frequency of the probe, i.e., $r_2 \approx \sqrt{\frac{4\pi}{K_B T}}$, with $D$, the vortex diffusion constant of the vortices. The main effect of $\epsilon_2$ is to induce a small tail above $T_{BKT}$ for the finite-frequency stiffness, as given by $J_f = \frac{1}{k_B T} \ln (\xi/\xi_0)$, as we indeed observe in the experiments. Here we follow the same procedure outlined in ref. \(^{38}\) to compute $\epsilon(\omega)$, and in full analogy, with this previous work, we find a very small vortex diffusion constant $D \approx 10^{-3} \text{nm}^2/\text{s}$. The correlation length $|t|$ also enters the temperature dependence of the resistivity above $T_{BKT}$, that follows the usual scaling $R/R_0 = 1/|t|$. To interpolate between the BKT and Gaussian regime of fluctuations, we use the well-known Halperin–Nelson expression \(^{41,42}\) for the BCS stiffness, as given by $\epsilon(\omega) = \frac{1}{4\sqrt{\pi}} \frac{\omega}{|k_B T|} \ln (\xi/\xi_0)$, with $\xi = \sqrt{\frac{4\pi}{\omega}}$.

Finally, to account for the possible inhomogeneity of the sample, we consider the extension of the previous method to the case where the overall complex conductivity of the sample is computed in the selfconsistent effective-medium approximation \(^{27}\) as the solution of the following equation:

$$\sum_i P_i \frac{\sigma_i(\omega) - \sigma(\omega)}{\sigma_i(\omega) + \sigma(\omega)} = 0 \quad (4)$$

Here $\sigma_i(\omega)$ denotes the complex conductivity of a local superconducting puddle with stiffness $J_f$ and local $T_f$, that are taken with a Gaussian distribution $P_i$ with variance $\sigma_c$, centered around the BCS fit of $f_{EXP}^{\exp}$. For each realization $J_f$ we then compute the $J_f$ from the solution of the BKT equations \(^{27}\).-\(^{38}\), we determine the corresponding complex conductivity $\sigma_i(\omega)$ and we finally solve Eq. (4) to get the average $f^{BKT} = (n^2/4\pi^2) \omega \sigma_i(\omega)$ below $T_f$, and the average $\sigma(\omega) = 1/R^{\exp}$ above $T_f$, i.e., the dashed lines reported in Fig. 4c. Further details about the implementation of the effective-medium approximation can be found in refs. \(^{16,38}\). The main effect of inhomogeneity is to contribute slightly to the suppression of $J_f$ with respect to $J_{BCS}$ before $T_{BKT}$. In our case, we checked that inhomogeneity, if present, is very small, and a $\sigma_c = 0.02$ is enough to account for the measured temperature dependences.

Data availability

The authors declare that the data that support the findings of this study are available within the article. All other relevant data are available from the corresponding authors upon request.

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Author contributions

N.B. and M.B. proposed and supervised the study. S.M. and H. W. prepared the samples and performed XPS experiments and their analysis. A.G. performed the STEM and EELS analysis. G.C.M., G.S., and S.M. performed the dc and microwave transport experiments and analysed the data with input from M.B., J.L., and N.B. L.B. conducted the BKT analysis of microwave data. N.B., M.B., and L.B. wrote the manuscript with input from all authors. All authors discussed the results and contributed to their interpretation.

Competing interests

The authors declare no competing interests.

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