Signatures of a strange metal in a bosonic system

Fermi liquid theory forms the basis for our understanding of the majority of metals: their resistivity arises from the scattering of well defined quasiparticles at a rate where, in the low-temperature limit, the inverse of the characteristic time scale is proportional to the square of the temperature. However, various quantum materials—notably high-temperature superconductors—exhibit strange-metallic behaviour with a linear scattering rate in temperature, deviating from this central paradigm. Here we show the unexpected signatures of strange metallicity in a bosonic system for which the quasiparticle concept does not apply. Our nanopatterned YBa$_2$Cu$_3$O$_7$ (YBCO) film arrays reveal linear-in-temperature and linear-in-magnetic-field resistance over extended temperature and magnetic field ranges. Notably, below the onset temperature at which Cooper pairs form, the low-field magnetoresistance oscillates with a period dictated by the superconducting flux quantum, $h/2e$ ($e$, electron charge; $h$, Planck’s constant). Simultaneously, the Hall coefficient drops and vanishes within the measurement resolution with decreasing temperature, indicating that Cooper pairs instead of single electrons dominate the transport process. Moreover, the characteristic time scale, $\tau$, in this bosonic system follows a scale-invariant relation without an intrinsic energy scale: $h/\tau = a(k_B T + \gamma \mu_B B)$, where $h$ is the reduced Planck’s constant, $a$ is of order unity, $k_B$ is Boltzmann’s constant, $T$ is temperature, $\mu_B$ is the Bohr magneton and $\gamma = 2$. By extending the reach of strange-metal phenomenology to a bosonic system, our results suggest that there is a fundamental principle governing their transport that transcends particle statistics.

The conventional theory of metallic transport predicts that the electrical resistivity $\rho(T)$ at low temperatures should vary as a function of $T^2$, owing to the dominance of electron–electron scattering. However, this common understanding has encountered formidable conceptual challenges in accounting for some recent discoveries. One emerged soon after the discovery of high-temperature superconductors, in the form of the linear temperature dependence of their normal-state resistivity. In sharp contrast to a quasiparticle picture. Interestingly, the high-temperature region naturally involves collective excitations due to the Fermi liquid theory forms the basis for our understanding of the majority of metals: their resistivity arises from the scattering of well defined quasiparticles at a rate where, in the low-temperature limit, the inverse of the characteristic time scale is proportional to the square of the temperature. However, various quantum materials—notably high-temperature superconductors—exhibit strange-metallic behaviour with a linear scattering rate in temperature, deviating from this central paradigm. Here we show the unexpected signatures of strange metallicity in a bosonic system for which the quasiparticle concept does not apply. Our nanopatterned YBa$_2$Cu$_3$O$_7$ (YBCO) film arrays reveal linear-in-temperature and linear-in-magnetic-field resistance over extended temperature and magnetic field ranges. Notably, below the onset temperature at which Cooper pairs form, the low-field magnetoresistance oscillates with a period dictated by the superconducting flux quantum, $h/2e$ ($e$, electron charge; $h$, Planck’s constant). Simultaneously, the Hall coefficient drops and vanishes within the measurement resolution with decreasing temperature, indicating that Cooper pairs instead of single electrons dominate the transport process. Moreover, the characteristic time scale, $\tau$, in this bosonic system follows a scale-invariant relation without an intrinsic energy scale: $h/\tau = a(k_B T + \gamma \mu_B B)$, where $h$ is the reduced Planck’s constant, $a$ is of order unity, $k_B$ is Boltzmann’s constant, $T$ is temperature, $\mu_B$ is the Bohr magneton and $\gamma = 2$. By extending the reach of strange-metal phenomenology to a bosonic system, our results suggest that there is a fundamental principle governing their transport that transcends particle statistics.

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resistance behaviour can occur in a bosonic host. The nanopatterned YBCO films exhibit a T-linear resistance, $h/2e^2 \times 1/T_c^{onset} (T_c^{onset})$, the onset temperature at which Cooper pairs form. This resembles the one in fermionic systems, $a_0 = h/2e^2 \times 1/T_c$. Moreover, the magnetoresistance is linear in this regime with a slope $\beta_{\psi_0}$ that appears comparable to $h/2e^2 \times 1/B^* (B^* \text{ is a characteristic magnetic field})$. Interestingly, the ratio of the slopes, $a_0/\beta_{\psi_0}$, appears to be determined only by fundamental constants, $2\mu_\psi/\hbar$, indicating the lack of an intrinsic energy scale of the system and therefore scale-invariant behaviour. Our study expands the possible quantum space for strange-metal states and advances the theoretical understanding of the underlying physics.

**T-linear resistance**

The present studies were conducted on a series of high-$T_c$ superconducting YBCO films that were patterned from pristine YBCO films by reactive ion etching (RIE) through a nano-porous mask with a triangular array of holes (Extended Data Fig. 1). Nanopatterning left these samples with a hexagonal array of superconducting ‘islands’ connected by weak links (Fig. 1b). Increasing the etching dose can increase the normal-state resistance $R_N$ by inducing more side wall damage in the pores to make the links between array nodes more resistive. As a result, our nanopatterned samples feature a range of $R_N$.

As shown by the sheet resistance as a function of temperature measurements, $R_s(T)$, in Fig. 1a, these arrays undergo a superconductor–anomalous metal–insulator transition with increasing $R_N$. The transport behaviour that we focus on occurs in films near the anomalous metal–insulator phase boundary where $R_s$ is close to $2R_N$, where $R_N$ is the quantum resistance, $h/4e^2$. For $R_N$ less than about 3 kΩ, the $R_s(T)$ exhibits transitions to a superconducting state starting at an onset temperature $T_c^{onset}$ that is close to the pristine film transition temperature. As $R_N$ grows larger than approximately 3 kΩ, the anomalous-metal phase for which $R_s(T)$ saturates at low temperatures to a constant value appears (see Supplementary Fig. 1). When $R_N$ approaches 12 kΩ, near the anomalous metal–insulator transition point, $R_s(T)$ develops a nearly linear temperature dependence over an extended range, as shown by six representative samples, f0–f5. In the case of f2 ($R_N$ approaches 12 kΩ), the truly linear temperature-dependent behaviour is observed from 8 K to 43 K (that is, $dR_s/dT = 160 \Omega K^{-1}$) with a low-temperature intercept near zero (Extended Data Figs. 2–4). Above 12 kΩ, the linear temperature dependence develops into an insulating dependence (that is, $dR_s/dT < 0$) in the low-temperature limit.

Having detected a regime with T-linear resistance, we now turn to investigate its slope. The extracted slope $\alpha_{\psi_0} = dR_s(T)/dT$ values for each sample measured in Fig. 1a are plotted against $R_N$ in Fig. 1c. Over the $R_N$ range that brackets f2, 10–15 kΩ, the slope remains within 25% of the maximum value, whereas the lowest temperature behaviours of the $R_s(T)$ evolve from anomalous metal (for example, f3 and f5) to insulator (for example, f0). Interestingly, $\alpha_{\psi_0}$ of the samples near 12 kΩ is close to $h/2e^2 \times 1/T_c^{onset}$. Then, we introduce a normalized parameter $\eta = (2e^2/h)\alpha_{\psi_0} T_c^{onset}$, and $\eta$ is close to 1 around the critical region, as

![Image 186x631 to 377x738]
shown in Fig. 1c. By contrast, the bounded slope in a fermionic strange metal is characterized by $\eta_\ell = (2e^2/h)\alpha_\ell T_c$, where $0.3 \leq \eta_\ell \leq 1.5$ (ref. 5). Notably, the slope $\alpha_\ell$ in bosonic systems is much larger than $\alpha_\ell$ in fermionic systems, because $T_c$ is usually two orders of magnitude larger than $T_c^{*}$.

**Bosonic nature**

Below $T_c^{*}$, two major characteristics indicate that Cooper pair motion dominates the transport in the $T$-linear and $B$-linear resistance regimes. (1) The magnetoresistances (Fig. 1d and Supplementary Fig. 2) oscillate with a constant period of $\mu_B H_{\text{period}} = 0.225 T$ (with vacuum permeability), corresponding to one superconducting flux quantum per unit cell of the array with $H_{\text{period}} = h/2eS$ and the unit cell area $S = 9,200 \, \text{nm}^2$. These $h/2e$ magnetoresistance oscillations provide direct evidence of Cooper pair transport$^{30,31}$. (2) The Hall coefficient $R_{xy}$ drops below $T_c^{*}$ to below our resolution at 30 K (Fig. 1e). The zero $R_{xy}$ reveals that the transport is particle–hole symmetric. Its overlap with the oscillations shows that the particle–hole symmetry reflects the dominance of Cooper pair transport. By contrast, the Hall coefficient $R_{xy}$ in fermionic systems (for example, LSCO and BaFe$_2$(As$_{1-x}$P$_x$)$_2$)$^{32,33}$ is finite and saturates with decreasing temperatures near the quantum critical point. The typical resolution for the Hall resistance of sample f4 (as shown in Fig. 2f, g) is 0.2 $\Omega$.

**B-linear resistance**

While there is magnetoresistance, the $T$-linear resistance persists in magnetic fields ($B > 1 \, \text{T}$) beyond the magnetoresistance oscillations, as shown in Fig. 2a, b. Below $T_c^{*}$, the magnetoresistance is uniformly positive and nearly linear up to the measurement limit of 9 T (Fig. 2c, d and Extended Data Figs. 5a–e, 6). By contrast, the magnetoresistance of a lower $T_c$, superconducting sample is nonlinear (Extended Data Fig. 5f). Meanwhile, the Hall coefficient $R_{xy}$ drops to near zero in most of the $B$-linear resistance regime up to 9 T (Fig. 2e).

Notably, the magnetoresistance curves of samples f0–f5 are approximately parallel for a wide range of temperatures. Their slopes obtained from linear fits, $\beta_{cp} = dR/dB$, as a function of temperature capture this observation (Fig. 2f, g and Supplementary Fig. 3). $\beta_{cp}$ for film f1, for example, shows a plateau over which it varies by only 12% from 30 K to 2 K, while its resistance changes by almost a factor of 3 as shown in Fig. 2g. The temperature range for this plateau or broad peak coincides with the temperature range over which the resistance is linearly temperature dependent. Writing the magnetoresistance in a form similar to the temperature dependence $\frac{dR}{dT} = \beta_{cp} T$, we can identify a magnetic field scale $B^*$ for the slope. At the maximum slope, $B^*$ is approximately 60 T, which is comparable to the upper critical field $B_{c2}$ of a YBCO film near optimal doping$^{32}$.

**Scale-invariant resistance**

The linear dependencies of the sheet resistance on $T$ and $B$ indicate that

$$
\frac{R_s(T, B) - R_s(0, 0)}{T} = \alpha_{cp} \frac{B}{T} \beta_{cp}
$$

(1)

over a wide temperature and magnetic field range. Replotting the $R_s(T, B)$ as $(R_s(T, B) - R_s(0, 0))/T$ versus $B/T$ leads to the collapse onto a single curve from 3 K to 30 K as shown for f1 in Fig. 3a. The scaling...
plot of other samples are shown in Fig. 3b and Extended Data Fig. 7. The ratio of $R_{\text{onset}}$ and $a_\text{on}$ gives a quantity that is close to a ratio of the two fundamental constants: $k_B/\hbar\nu_\text{b}$, where $\nu_\text{b}$ is close to 2 (Extended Data Fig. 8). Both $T$-linear and $B$-linear resistance imply scale-invariant damping rates that satisfy $\hbar/\tau = a \kappa T$ and $\hbar/\tau = a \eta \mu_B B$, respectively ($a$ is of order unity).

**Phase diagram**

To depict the region over which the scaling appears to hold, we plot $1/R_{\text{onset}} \times d^2 R/ dT^2$ on a colour scale on $R_{\text{onset}}$ versus $T$ axes, as shown in Fig. 4a. This region is represented by the wide white region in the central part Fig. 4a, where $1/R_{\text{onset}} \times d^2 R/ dT^2$ is near zero. It funnels down with decreasing temperature from $T_{\text{onset}}$ becoming narrow near the critical sheet resistance $R_{\text{wc}} = \hbar/2e^2$, separating the insulating and anomalous-metal phases at low temperatures. We label this white area as the bosonic strange metal, reminiscent of the strange metal in fermionic systems.

**Discussion**

The present work uncovers a distinctive low-temperature bosonic metallic transport behaviour that deviates from the previous boson localization paradigm of the superconductor–insulator transition in which metallic behaviour appears with a critical resistance $R_c$ at a critical point of a tuning-parameter–like magnetic field or normal-state resistance. Similar to the anomalous-metal state, this strange-metal state exists over a range of tuning parameters with residual resistances far below $R_c$. The primary characteristics that distinguish this bosonic transport behaviour are: (1) a resistance that changes linearly with temperature, with a slope that appears comparable to a ratio of a resistance quantum and an energy scale characterizing the ground state of the charge carriers; (2) a magnetoresistance that changes linearly with magnetic field as implied by the data collapse in Fig. 3; and (3) the appearance of these behaviours in a fan-shaped region that appears to terminate in a quantum critical point in the phase diagram in Fig. 4.

Because these bosonic transport characteristics share some features with fermionic strange metals, we discuss them in that context first. In the model based on the Drude approach, the scattering rate in fermionic systems with Planckian dissipation follows $1/\tau = a k_B T/\hbar$, where $a$ is of order unity. Correspondingly, the slope $dR_c(T)/dT$ of the sheet resistance, $a = (m^*/n^*) (k_B/\hbar)$, then depends only on $m^*/n$, where $m^*$ is the effective mass and $n$ is the carrier density. For fermionic systems, $n/m^*$ is proportional to $k_B T_c$ in a two-dimensional approximation. Thus $a = \alpha/2\pi \approx 1/\gamma^2$ applies for fermionic systems. By analogy, in the nanopatterned two-dimensional superconducting array, the bosonic carriers dominate the transport process and correspondingly, $n_\text{on}/m^*$ is proportional to the characteristic energy scale $k_B T_{\text{onset}}^*$, where $n_\text{on}$ is the density of Cooper pairs on an island. Consequently, $a_\text{on}$ of this bosonic system is much larger than that of fermionic systems ($n_\text{on}/n_c = T_{\text{onset}}^*/T_{\text{c onset}}^*$), consistent with the transport data (also see Supplementary Table 1). The Hall measurements reveal an important qualitative difference between the bosonic and fermionic systems. For the bosonic strange metal (Fig. 1e), $R_{\text{onset}}$ drops to zero with decreasing temperature below $T_{\text{c onset}}^*$, whereas for fermionic strange metals, $R_{\text{onset}}$ saturates to a finite value in the zero-temperature-limit (for example, LSCO and BaFe$_2$(As$_{1-x}$P$_x$)$_2$). To conclude, the Hall signal provides a qualitative difference between fermionic and bosonic strange metals.

Within the context of the natural model for these arrays, the quantum XY model with dissipation provides an explanation for why this behaviour has not been observed before. The dissipative processes involve the coupling of the bosons to low-lying fermionic excitations of this d-wave system. This dissipative model can exhibit three ground states rather than two for non-dissipative models. Namely, as shown in Fig. 4b, it has an ordered superconducting state, a disordered insulating state and, in between, a quasi-ordered state that is of interest here. The phase of the order parameter in the less familiar quasi-ordered state is spatially ordered and temporally disordered (Fig. 4c). Here, by increasing the normal-state resistance $R_{\text{onset}}$, one crosses through these three states, as shown in Fig. 4b. The absence of strong-metal behaviour in conventional superconducting systems that exhibit anomalous-metal behaviour suggests that its emergence here is related to the high-$T_c$ superconducting state. In previous investigations, their anomalous-metal behaviours appear at substantially lower $R_{\text{onset}}$, where the Berezinskii–Kosterlitz–Thouless (BKT) transition is closer to $T_{\text{c onset}}^*$ (ref. 23). Moreover, the higher $T_{\text{c onset}}^*$ in nanopatterned YBCO films can make their linear- $T$ resistance more apparent as it extends over a wide range of temperature.
metallic state and its underlying physics. Also, future measurements of the frequency-dependent conductivity show promise for direct insights into the dynamics of this system, which could reveal the origins of these linear transport phenomena. A measurement of the Drude width would be particularly interesting, as it has been proposed to scale with the entropy in models applicable to matter in many forms including cold atoms, solid-state materials, and quark gluon plasma. In addition, as dissipation-induced decoherence processes are important in quantum computing, our results provide a platform for further study of the microscopic origins of the dissipative strange-metal state.

Online content

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**Methods**

**Sample preparation**

12-nm-thick high-Tc YBa2Cu3O7−δ superconducting films were epitaxially grown on (001) oriented substrates SrTiO3 in d.c. magnetron sputtering system using an argon/oxygen (ratio 2:1) mixture under a working pressure of 30 Pa at T = 700 °C, with a growth rate of ~0.5 nm min−1. The oxygen stoichiometry (δ = 0.1, p = 0.16) was determined by comparing the transport and Φc of our sample with previously reported YBCO films with varying doping level31, as shown in Supplementary Fig. 4. Then, an anodic aluminium oxide (AAO) template 200-nm thick was transferred to YBCO thin film. By reactive ion etching, the AAO pattern of a triangular array of holes with a 70-nm diameter and 103-nm period is duplicated onto the YBCO films. The etching recipe includes 30 sccm CHF3, 40 sccm Ar, 20 mtorr (1 torr ≈ 133 Pa) pressure, 200 W RF power, 1,800 W inductive coupling plasma (ICP). The YBCO films can be tuned from a superconducting state to an insulating state by increasing the etching time from 20 s to 160 s30. By increasing the first anodization time we can prepare YBCO films with a more periodic array of honeycomb holes. As a result, the samples in this study show more uniform link resistance than in a previous study30.

**Measurements**

The in-plane resistance and magnetoresistance were measured with a standard four-probe in a commercial Physical Property Measurement System (PPMS, Quantum Design) with magnetic fields up to 9 T. For the transport measurements, ten distinct samples were measured (f0−f9) (see Supplementary Table 2). For standard four-probe measurements, the channel size of the YBCO sample is 1 mm × 1 mm, as shown in Extended Data Fig. 9a. For Hall measurements, channel size of the YBCO sample is 1 mm (length) × 0.9 mm (width), as shown in Extended Data Fig. 9b. Low excitation currents (10 nA−1 μA) were applied in the ab plane and the magnetic field was applied along the c axis. As shown in Extended Data Fig. 10, the current voltage (I−V) curves are linear and the R−T curves are independent of measuring current over two orders of magnitude, indicating that the measurements were in the ohmic regime. The sweeping rate is typically 0.04 T min−1 for measuring the magnetoresistance quantum oscillations.

**Comparison of the T-linear and B-linear resistivity**

The T-linear resistivity has been detected in various fermionic systems42–43. The strange-metal phase has received considerable attention, owing to its potential relevance to so-called Planckian dissipation46–51. The appearance of the bosonic anomalous metallic state in a variety of two-dimensional systems52–56 indicates that the strange-metal phase may not be restricted to fermionic systems. Here we show the presence of this unusual T-linear resistance behaviour in a bosonic system, namely nanopatterned YBCO thin films. There are several prominent differences between the T-linear and B-linear resistivity of nanopatterned YBCO thin films and that of non-patterned LSCO30.

1. Below Tc,onset, the typical slope of the T-linear resistivity of nanopatterned YBCO films is two orders of magnitude larger than that for non-patterned LSCO (see Supplementary Fig. 5). The maximum slope of B-linear resistivity in nanopatterned YBCO thin film is also 2−3 orders of magnitude larger. As shown in Supplementary Fig. 6, the slope is 172 μΩ cm T−1 for sample f5, whereas it is only 0.32 μΩ cm T−1 for non-patterned LSCO. By contrast, above Tc,onset, the slope of T-linear resistivity for the non-patterned YBCO is comparable to that observed in non-patterned LSCO (see Supplementary Fig. 5, e). Moreover, the slope of the T-linear resistivity in nanopatterned YBCO thin films is enhanced by a factor of 10 when the temperature decreases below Tc,onset. For instance, the slope for f3 above Tc,onset is 12.9 μΩ cm K−1, which becomes around 180 μΩ cm K−1 below Tc,onset (see Supplementary Fig. 5).

2. The strange metal in nanopatterned YBCO thin films originates from bosonic carriers, which is verified by a vanishing Hall coefficient with decreasing temperature below Tc,onset. In the case of fermionic strange metal of non-patterned LSCO30, the Rn saturates to a finite value. As shown in Fig. 2e, the Hall resistance Rn is gradually suppressed and drops to zero with decreasing temperature below Tc,onset. Moreover, Fig. 1e shows that T-linear resistance Rn is extended to 60 K, and correspondingly the Hall coefficient Rn at 50 K (123 nΩ cm T−1) is already much smaller than that of the normal state (908 nΩ cm T−1) at 80 K. Finally, Rn drops to zero within the measurement resolution at a relatively high temperature of 30 K. Thus, the superconducting fluctuation effects may stabilize the T-linear resistance behaviour (from 50 K to 60 K), which leads to a slightly extended temperature regime for T-linear Rn, compared to the range with vanishing Rn30.

3. Next, we compare the characteristic time scale (τ) in both fermionic and bosonic strange metals. The scattering rate in fermionic systems follows either h/τ = (kBTc)3/2 (in BaFe2(As1−xP)x)14,58, overdoped Bi2Sr2CaCu2O8+δ (Bi2201)39 or h/τ = α(kcT + μcB) in LSCO30, where γ = 1 and α is of order unity. Here, in nanopatterned YBCO thin film, the characteristic damping rate follows a scale-invariant form: h/τ = α(kcT + μcB), where γ = 2 and α is of order unity.

**Impact of irradiation method**

Here we compare our results with what has been reported on non-nanopatterned YBCO films exposed to electron irradiation61. (1) The present nanopatterning produces superconducting ‘islands’ connected by weak links. This patterning leads to boson-dominated transport in the T-linear resistance regime, as has been confirmed by both h/2e oscillations and Hall resistance measurements. Electron irradiation, by contrast, depresses the order parameter throughout to reduce Tc (ref. 15). (2) The slope of T-linear resistance below Tc,onset in our samples is two orders of magnitude larger than what has been reported61. Thus, the bosonic nature and its T-linear resistive SIT is quantitatively analysed by the proliferation of mobile vortices above the BKT transition temperature21, whereas in 2D dissipative quantum phase transitions the resistivity21, whereas in 2D dissipative quantum phase transitions the resistivity may not be restricted to fermionic systems. Here we show the presence of this unusual T-linear resistance behaviour in a bosonic system, namely nanopatterned YBCO thin films. There are several prominent differences between the T-linear and B-linear resistivity of nanopatterned YBCO thin films and that of non-patterned LSCO30.

**Comparison with SIT in conventional and high-Tc superconductors**

In two-dimensional (2D) conventional superconducting systems, the superconductor–insulator transition (SIT) is always accompanied by a sharp decline around Tc,onset, owing to the BKT transition, which can be quantitatively analysed by the proliferation of mobile vortices above the BKT transition temperature52. Meanwhile, in high-Tc superconductors, the reported SIT shares similar features53,54 with conventional superconducting systems. To our knowledge, T-linear resistance has not been reported in SIT of non-nanopatterned films. The nanopatterned YBCO film arrays in this work are in a high-resistance limit, whereas previous investigations of cuprate arrays barely reach this regime.

Theoretically, the bosonic strange-metal phase in our observations may be related to the dissipative quantum phase transition tuned by the link resistance of a Josephson junction55. This dissipative quantum phase transition is controlled by the dimensionless parameter a = Rg/Rm, where Rm = h/4e2 denotes the quantum resistance and Rg stands for the local link resistance of the Josephson junction. The dissipative SIT occurs at certain critical value a, where a = 1 for 2D Josephson junction arrays with quasiparticle damping56. This dissipation-driven phase transition differs from disorder-tuned SIT mainly in two aspects. First, the local link resistance Rg needs to be homogenous throughout the sample. By contrast, a disorder-tuned SIT does not need to be in a homogenous system. Secondly, a 2D disorder-tuned SIT is always accompanied by a sharp resistance decrease around the BKT temperature52, whereas in 2D dissipative quantum phase transitions the resistance gradually declines with decreasing temperature. Although certain microscopic analysis has been conducted for the zero-temperature
regime of the anomalous-metal phase\textsuperscript{67}, a unified explanation for both bosonic strange-metal and anomalous-metal phases may need further investigation, which should capture the subtle damping process of the phase mode.

In our experimental setup, the etching time controls the link resistance $R_{\text{LN}}$ or the normal state resistance $R_{\text{s}}$ of the system and leads to the quantum phase superconductor–anomalous metal–insulator transition. We speculate that both bosonic strange metals and anomalous metals originate from the dissipation process, which is controlled by the link resistance $R_{\text{LN}}$ in the Josephson junction array\textsuperscript{21}. Subsequently, in high-$T_c$ superconducting films with gate-voltage-tuned SIT\textsuperscript{64,65}, the system cannot be considered as a Josephson junction array, and the absence of dissipation may prevent the appearance of the bosonic strange-metal phase.

Data availability

The data that support the plots within this paper are available from the Zenodo data repository, https://doi.org/10.5281/zenodo.5603259. Source data are provided with this paper.
Extended Data Fig. 1 | Scanning electron microscopy image of a nanopatterned $\text{YBa}_2\text{Cu}_3\text{O}_{7-\delta}$ (YBCO) thin film. The 12-nm-thick nanopatterned YBCO thin film was fabricated by reactive ion etching through an anodic aluminium oxide (AAO) membrane directly placed atop the YBCO. By RIE, the anodized aluminium oxide pattern of a triangular array of holes with ~70-nm diameter and ~103-nm period was duplicated onto the YBCO film.
Extended Data Fig. 2 | First derivatives of the R-T curves of nanopatterned YBCO films. a−f, The first derivatives of resistance as a function of temperature for f2 (a), f3 (b), f4 (c), f7(d) and f8 (e). The table shows the yielded parameters of the statistical analysis (f).
Extended Data Fig. 3 | Data residuals after subtracting the linear fit for the $R$–$T$ curves of nanopatterned YBCO films. Residuals for the $R$–$T$ curves of $f_2$ (a), $f_3$ (b), $f_4$ (c), $f_7$ (d) and $f_8$ (e). The residual is defined by the resistance subtracting the linear fitting of the $R$–$T$ curves with the slopes and intercepts shown in f. To delineate the temperature regime for $T$-linear resistivity, the residual cut-off is set by 50 $\Omega$ which is around 0.5% of the normal-state sheet resistance $R_N$. The table shows the temperature regime for $T$-linear resistivity where the residual is within 50 $\Omega$ (f).
Extended Data Fig. 4 | Nonlinear fitting for the $R$-$T$ curves of nanopatterned YBCO films. a–d. Least-squares nonlinear fitting of the $R$-$T$ curves for f8 (a), f7 (b), f4 (c) and f2 (d). $n$ is the yielded power from the fitting.
Extended Data Fig. 5 | Scale-invariant $\beta$-linear resistance in nanopatterned YBCO thin films under perpendicular magnetic field. a–f, The magnetoresistance for films: $f_0$ (a), $f_2$ (b), $f_3$ (c), $f_5$ (d), $f_6$ (e) and superconducting (SC; f).
Extended Data Fig. 6 | First derivatives and nonlinear fitting of the $R$–$B$ curves of nanopatterned YBCO films. 

**a, b**, The first derivative of resistance as a function of magnetic field for $f_4$ (**a**) and $f_2$ (**b**) at various temperatures. 

**c, d**, Least-squares nonlinear curve fitting of the $R$–$B$ curves for $f_4$ (**c**) and $f_2$ (**d**).
Extended Data Fig. 7 | $B$-$T$ scaling in nanopatterned YBCO films. a-c, $B$-$T$ scaling in nanopatterned YBCO thin films of f2 (a), f3 (b) and f5 (c).
Extended Data Fig. 8 | Magnetotransport as a function of $(k_B T + γμ_B B)/k_B$ of nanopatterned YBCO films. a–d, The resistance and magnetoresistance of $f_1$ (a), $f_2$ (b), $f_3$ (c) and $f_5$ (d) as a function of $(k_B T + γμ_B B)/k_B$, where the $γ$ parameter can be estimated by adjusting it when the curves collapse best.
Extended Data Fig. 9 | Electrode pattern for the measurement.

a, Illustration of the electrode pattern for standard four-probe measurements. The current was applied at electrode #1 and #5. The Hall resistance was measured from electrode #3 and #7, the longitudinal resistance is measured from electrode #2 and #3. STO, SrTiO$_3$.

b, Illustration of the electrode pattern for Hall measurements. The current was applied at electrode #1 and #5. The Hall resistance was measured from electrode #3 and #7, the longitudinal resistance is measured from electrode #2 and #3. STO, SrTiO$_3$. 
Extended Data Fig. 10 | Current voltage ($I$–$V$) curves and $R$–$T$ curves at different current excitations in nanopatterned YBCO film. a–c, $R$–$T$ curves for representative films f4 (a), f3 (b) and f2 (c) with different currents. d, Current voltage ($I$–$V$) curves for f4.