Quantum critical behaviour in magic-angle twisted bilayer graphene

Alexandre Jaoui1✉, Ipsita Das1, Giorgio Di Battista1, Jaime Díez-Mérida1, Xiaobo Lu1, Kenji Watanabe1,2, Takashi Taniguchi2, Hiroaki Ishizuka3,4, Leonid Levitov4 and Dmitri K. Efetov1✉

The flat bands\(^3\) of magic-angle twisted bilayer graphene (MATBG) host strongly correlated electronic phases such as correlated insulators\(^{24,25}\), superconductors\(^{26,27}\) and a strange metal state\(^{11}\). The strange metal state, believed to be key for understanding the electronic properties of MATBG, is obscured by various phase transitions and so it could not be unequivocally differentiated from a metal undergoing frequent electron–phonon collisions\(^{11}\). Here we report transport measurements in superconducting MATBG in which the correlated insulator states are suppressed by screening. The uninterrupted metallic ground state shows resistivity that is linear in temperature over three orders of magnitude and spans a broad range of doping, including that where a correlation-driven Fermi surface reconstruction occurs. This strange metal behaviour is distinguished by Planckian scattering rates and a linear magnetoresistivity. By contrast, near charge neutrality or a fully filled flat band, as well as for devices twisted away from the magic angle, we observe the archetypal Fermi-liquid behaviour. Our measurements demonstrate the existence of a quantum-critical phase whose fluctuations dominate the metallic ground state throughout a continuum of doping. Further, we observe a transition to the strange metal upon suppression of the superconducting order, suggesting a relationship between quantum fluctuations and superconductivity in MATBG.

Low-Fermi-energy metals, in which electron interactions are described by Landau’s Fermi-liquid theory, exhibit a resistivity $\rho$ that scales with temperature as $\rho \propto T^2$ and with magnetic field as $\rho \propto B^2$ (ref. $^{15}$). However, a different behaviour arises in the presence of strong electron correlations, where alongside superconductors, magnets and insulators, unusual strange metal phases can survive down to $T \to 0$, as reported in cuprates\(^{15,16}\), ruthenates\(^{17}\), pnicotides\(^{18}\), heavy-fermion systems\(^{19}\) and, recently, in twisted dichalcogenides\(^{20}\). Such strange metal phases display the unique dependences $\rho \propto T$ and $\rho \propto B$ and are associated with ultra-fast carrier scattering governed by the universal Planckian dissipation rate $1/\tau = k_B \theta / \hbar$, with $\hbar$ the reduced Planck’s constant and $k_B$ the Boltzmann constant\(^{22}\). The observation of $\rho \propto T$ down to $T \to 0$ signals the proximity to a quantum critical point: the neighbouring metallic ground state is dominated by inelastic scattering with electronic quantum-critical fluctuations. These fluctuations have been mostly related to magnetic ordering\(^{22}\), yet purely nematic fluctuations have also been reported\(^{23}\). Understanding the relationship between quantum fluctuations, finite-temperature strange metallicity and the exotic phase transitions found in a wide variety of strongly correlated systems is a major conundrum in condensed matter physics.

Early measurements on twisted bilayer graphene, in which the two layers are rotated by the magic angle of $\theta = 1.1^\circ$ (magic-angle twisted bilayer graphene (MATBG)), observed a metallic phase at temperatures $0.5 \, K < T < 30 \, K$ with a linear $\rho \propto T$ (refs. $^{12,13}$) and Planckian scattering rates in close proximity to correlated insulating\(^{24}\) and superconducting\(^{25}\) phases for electron filling factors of $n = \pm 2$ per moiré unit cell. First interpretations suggested a strange metal phase\(^{12}\), emerging from purely electronic interaction\(^{26}\) and drawing an analogy to other strongly correlated systems\(^{15–23}\). However, subsequently it was proposed that a conventional electron–acoustic phonon scattering mechanism\(^{27}\) above a characteristic Bloch–Grüneisen temperature $T_{BG}$ (typically, exceeding a few kelvin) could also result in a similar behaviour under favourable assumptions\(^{12,13}\). An investigation of the $T \to 0$ regime of MATBG would allow definitive differentiation between these distinct scenarios, yet, such studies have so far been impeded by the abundant low-temperature insulating and superconducting phase transitions.

In this Letter, we report on comprehensive electronic transport measurements extending down to unprecedentedly low temperatures of 40 mK, firmly establishing the existence of a strange metal phase. In order to reveal the low-temperature metallic states, we deliberately chose MATBG devices (listed in Supplementary Section A) with ultra-close metallic screening layers and twist angles that slightly deviate from $1.1^\circ$ (refs. $^{9,27}$). In particular, we focus on device D1 with a twist angle of $\theta = 1.04^\circ$ and a screening layer spacing of $d = 9.5 \, nm$ (inset of Fig. 1a), which did not show correlated insulating phases (in the hole-doped region) but still showed a robust superconducting dome. We performed resistivity measurements for a broad parameter space of $(n,T,B)$ with carrier density $n$ tuned across the entire flat-band region, temperatures from $T = 40 \, mK$ to $T = 20 \, K$ and magnetic fields up to $B = 1 \, T$ (see Supplementary Section B for extended measurements up to 100 K). Our measurements confirm the existence of a $T$-linear resistivity in the centre of the flat bands above $T > 1 \, K$, in agreement with previous results\(^{12,13}\) and firmly establish its uninterrupted continuation down to 40 mK, which cannot be explained by conventional electron–phonon collisions. Further, we unveil another signature of the strange metal state: a linear $B$-dependence of the resistivity. By contrast, we find a typical Fermi-liquid behaviour $\rho \propto (T^2B^2)$ in the vicinity of the band edges and in twisted bilayer graphene with twist angles that strongly deviate from the magic angle $\theta > 1.3^\circ$ (see Supplementary Section C for more details). Lastly, we demonstrate that the strange
metal state extends into the superconducting dome region after suppressing superconductivity by an applied magnetic field. These observations establish the existence of a quantum-critical continuum and demonstrate a strong resemblance between MATBG and a variety of quantum-critical systems, pointing to an intrinsic relationship between electronic quantum fluctuations and superconducting/correlated insulating states. While recent findings of the isospin Pomeranchuk effect in MATBG28,29 point to soft spin and valley fluctuations as a driving mechanism of the low-$T$ phase diagram, the nature of the quantum fluctuation in MATBG remains an open question.

Figure 1a shows the four-terminal resistivity $\rho$ of device D1 as a function of the moiré band filling factor $\nu = n/n_0$, with $n_0 = 1/A_0$ ($A_0$ is the area of the moiré unit cell) for temperatures ranging from $T = 40$ mK up to $T = 20$ K. We observe insulating behaviour, that is an increasing resistivity when temperature is decreased, at electrostatic doping levels which correspond to the vicinity of the charge neutrality point $\nu = 0$, for the fully filled flat band $\nu = \pm 4$ and around $\nu = \pm 3$. Superconducting domes are found close to half filling ($\nu = \pm 2$). By choosing a screening layer separation that is smaller than a typical Wannier orbital size of 15 nm (ref. 9) we were able to quench the correlated insulators at $\nu = \pm 2$ and leave the hole-doped region entirely metallic (apart from the superconducting dome). The isospin Pomeranchuk effect is also clearly resolved, where especially the resistance peaks at $\nu = \pm 1$ are more pronounced at elevated $T$ (refs. 28,29). The simplicity of the phase diagram of the hole-doped D1 allows for an in-depth study of the metallic ground state and its ties to the neighbouring states. Similar datasets measured on other devices, shown in Supplementary Sections B and C, draw a consistent picture of the metallic ground state.

We first discuss the temperature dependence of the resistivity and its evolution with the filling factor. A basic zero-order picture emerges from the (numerical) derivative $(d\rho/dT)$, shown in Fig. 1b: this quantity is roughly temperature independent over a wide range of filling factors and is weakly sensitive to doping. The details of this behaviour are illustrated in Fig. 1c which presents the resistivity versus temperature for successive filling factors. Starting from the insulating regime at the charge neutrality point, metallicity is recovered at $\nu \approx -0.15$, which first shows a super-linear temperature dependence below $T < 15$ K and then saturates into a linear dependence. With increased doping, the onset of the linear dependence is quickly shifted to lower temperatures. Starting from $\nu \approx 2$, the $T$-linear regime extends down to the base temperature and remains $T$-linear until a second super-linear regime is found for $\nu < -3.5$. This strange metal phase is only interrupted by a superconducting transition observed around half filling.

We analyse the temperature dependence below $T < 10$ K by fitting the resistivity with $\rho(T) = \rho_0 + A_2 T^\gamma$, where the parameters $A_2$ and $\gamma$ define the prefactor and the exponent of the power law $T$-dependence, respectively and $\rho_0$ is the residual resistance at $T = 0$. Figure 1d shows $(\rho(T) - \rho_0)$ on a log–log scale, which allows $\rho(T)$ to be traced over more than three orders of magnitude down to centikelvin temperatures. We fit each curve with $A_2 T^\gamma$, which results in linear lines on the log–log plot, with a slope that is directly defined by $\gamma = d\ln[\rho(T) - \rho_0]/d\ln(T)$. In the proximity of the charge neutrality point ($\nu \approx -0.2$) and of full filling ($\nu = -3.7$), we find a $\gamma = 2 \pm 0.1$ dependence which results in a super-linear $\rho(T) \propto T^{2.0 \pm 1}$ dependence. However, for the filling factor range of $-3.5 < \nu < -2$ we find $\gamma = 1$, which gives rise to a strictly linear $\rho(T) \propto T$ dependence with an ultra-high filling-dependent slope of $A_2 > 0.25 \Omega$ K$^{-1}$ (shown in the inset of Fig. 1d). Strikingly, the $T$-linear dependence extends without interruption from the base temperature of 40 mK to a temperature of 10 K, above which it saturates (see Supplementary Sections D and E for a discussion on this saturation upon reaching the Mott–Joffe–Regel limit). The unique aspects of this ’zero-temperature’ transition are discussed in detail below.

Can the observed linear dependence be explained by electron–phonon scattering? The electron–phonon mechanism yields a weak $T$-linear resistivity only above the Bloch–Grüneisen temperature with typical values $T_{BG} > 10$ K, whereas below $T_{BG}$ the $T$-dependence
is super-linear, $\rho_{ph} \propto T^4$ (ref. 25). The temperature of 40 mK at which the $T$-linearity is observed in device D1 is three orders of magnitude lower than $T_B$, and the observed slope of $A_{ij} \approx 0.25$ kΩ K$^{-1}$ is much higher than expected from an electron–phonon mechanism. Since $T_B$ is proportional to the square root of $n$, $T_B \propto \sqrt{n}$ (ref. 25), it has been suggested that near the charge neutrality point ($\nu = 0$) or the Fermi energy resets at ($\nu = \pm 2$), the values $n$ and $T_B$ become small yielding a $\rho_{ph} \propto T$ dependence with an enhanced slope that can persist to temperatures as low as $T = 0.5$ K (refs. 13,26). However, this scenario is inconsistent with our findings as we observe (1) a large interval of dopings where the low-$T$ linear regime occurs, (2) the emergence of a $T$–dependent resistivity of the same amplitude near the flat-band edges and (3) the evolution of the prefactor $A_{ij}$ as a function of doping, which sharply increases for $\nu < -3$, as shown in the inset of Fig. 1d.

These low-temperature observations are summarized in a schematic phase diagram in Fig. 2a and are overall better explained by a purely electronic phenomenon. The metallic ground state at the edges of the band is a Fermi liquid that displays a quadratic $T$–dependent resistivity. However, doping away from the band edges induces two zero-temperature phase transitions, to a non-Fermi-liquid state with a $T$–linear resistivity for filling factors $-3.5 < \nu < -2$. Thus, it is tempting to identify the latter state as a strange metal above a quantum-critical phase, wherein the finite-$T$ metallic properties are dominated by critical fluctuations15–20. Support for this interpretation is offered by a comparison to the celebrated property of strange metals—a quasiparticle scattering time (ref. 25), a rough estimate of the scattering rate (Supplementary Section D) indicates that the inelastic $T$–linear resistivity is indeed consistent with Planckian dissipation, confirming previous reports12. Overall these findings resemble the $T$–linear resistivity observed in La$_2$Sr$_2$CuO$_4$ (LSCO) and Bi$_2$Sr$_2$CaCu$_2$O$_{8+}$ (Bi2201) (ref. 31) down to $T \rightarrow 0$, both at and away from a quantum critical point and are possibly associated with the reconstruction of the Fermi surface15.

Additional evidence for the quantum-critical continuum is provided by the superconducting state close to $\nu \approx -2$, where we
In a wide variety of strongly correlated systems, in spite of the band edges, a saturation is seen at low field \( B < 0.2 \, \text{T} \), which marks the crossover between thermal and magnetic dominated scattering times \( \tau_1 = \tau_2 \). The inset shows the evolution of the linear and quadratic prefactors \( A_B \) and \( \gamma \) versus \( T \), slopes that interact with the metallic ground state, which are characteristic of the superconducting-suppressed state from the zero-field magnetic properties.

Overall we find a behaviour very similar to the temperature dependence. In the proximity of the charge neutrality point (\( \nu = -0.2 \)), for \( T < T_c \), as shown in Fig. 2c for 40 mK, above \( B_c \), the magnetoresistivity (MR) scales linearly as \( \rho(B) \propto B \) up to \( B = 1.5 \, \text{T} \). This allows us to evaluate the MR-corrected resistivity, where we assume a \( -1 \) dependence, which gives rise to the power-law scaling of the MR. We illustrate the evolution of the MR across the band, where in Fig. 3a we observe a recovery of the strange metal phase upon suppression of the superconducting dome by a small perpendicular magnetic field \( B_c = 300 \, \text{mT} \) (Fig. 2b).

In addition, the \( T \)-linear resistivity is also accompanied by a \( B \)-linear MR outside the superconducting dome. Such linear MR provides additional evidence for the existence of critical fluctuations that interact with the metallic ground state, which are characteristic of strange metal phases in a multitude of superconductors. We find a behaviour very similar to the temperature dependence. In the proximity of the charge neutrality point (\( \nu = -0.2 \)), for \( T < T_c \), as shown in Fig. 2c for 40 mK, above \( B_c \), the magnetoresistivity (MR) scales linearly as \( \rho(B) \propto B \) up to \( B = 1.5 \, \text{T} \). This allows us to evaluate the MR-corrected resistivity, where we assume a \( -1 \) dependence, which gives rise to the power-law scaling of the MR.
to a linear $\rho \propto B$ dependence with an ultra-high filling-dependent slope of $\rho_B > 2k_\text{B} \text{K}^{-1}$ for $B > 0.2T$ (shown in the inset of Fig. 1d), but which saturates for $B \rightarrow 0$. We additionally show similar concomitant $B$-linear (up to $3T$) and $T$-linear resistivities in another device with nearly identical twist angle (only 0.01° apart) in Supplementary Section H. A similar study of the electron-doped region, although made more difficult because of the presence of successive phase transitions, also highlights the emergence of a $(B,T)$-linear resistivity in the centre of the electron-doped flat band (Supplementary Section I).

The observed linear MR is distinct from previous reports of linear MR for graphene systems. In these reports, the linear MR originates from classical effects and persists to $T = 300 \text{K}$ and $B = 62 \text{T}$ (ref. 35). Our MR is approximately 100 times stronger and is considerably more fragile, as it saturates near 1T and is quickly suppressed at elevated $T$, as is illustrated in Fig. 3d. As can be seen in Fig. 3e, between 40mK and 50K the slope of the MR $(\rho_B)$ decreases tenfold. By contrast, for devices with $\theta > 1.3\degree$, we observe an almost absent MR (Supplementary Section C). The identical scaling for the $B$- and $T$-dependences suggests the absence of an intrinsic energy scale in the ground state. We propose that it is dominated by scattering off quantum fluctuations, with a quasiparticle scattering rate which is given by the dominant energy scale $h/\tau = \max \{ h/\tau_m, h/\tau_T\} = \max \{ \beta \mu_\text{B} B, \alpha k_\text{B} T\}$, where $\tau_m$ and $\tau_T$ are the magnetic and thermal scattering times, $\mu_\text{B}$ is the Bohr magneton and $\alpha$ and $\beta$ are numerical factors. We estimate that at $1.9 \text{K}$, $\tau_m/\tau_T > 0.25 \text{T}$ (see Supplementary Section P for a detailed discussion), which corresponds exactly to the onset of the saturation observed in the low-field MR in Fig. 3e. Hence, for $B \rightarrow 0$ we can conclude that the MR saturates because of finite temperature effects, similar to the case of LSCO 36. On the contrary, the ansatz proposed to describe transport relaxation rates of pnictides, cuprates, heavy fermions and twisted chalcogenides near a quantum critical point, $h/\tau = (\alpha k_\text{B} T)^2 + (\beta \mu_\text{B} B)^2$ (refs. 36, 37), does not account for the MR of MATBG at finite temperatures.

These findings make a clear case that MATBG possesses a Planckian-limited $T$-linear resistivity that extends down to unprecedentedly low temperatures of 40mK and occurs alongside a quantum $B$-linear magnetoresistance. Such behaviour is incompatible with a Fermi-liquid picture and conventional electron–phonon scattering. The Fermi-liquid behaviour is observed throughout the entire moiré band at non-magic angles (as shown in a study of a variety of twist angles discussed in Supplementary Section B); by contrast, in MATBG it is pushed to the flat-band edges. We therefore conclude that a strange metal phase exists, arising from a quantum-critical region spanning a range of dopings including but not limited to those where the Fermi surface reconstructs and where quantum fluctuations dominate the metallic ground state of MATBG. Our observations reveal that superconductivity in MATBG emerges in a state dominated by quantum fluctuations. While the precise relationship between quantum fluctuations and phase transitions, as well as the microscopic nature of the fluctuations, remains an open problem, our work establishes a clear connection between strongly correlated and highly tunable electronic moiré systems and the universality class of the quantum critical matter.

**Online content**

Any methods, additional references, Nature Research reporting summaries, source data, extended data, supplementary information, acknowledgements, peer review information; details of author contributions and competing interests; and statements of data and code availability are available at https://doi.org/10.1038/s41567-022-01556-5.

Received: 27 July 2021; Accepted: 18 February 2022; Published online: 11 April 2022

**References**

1. Bistritzer, R. & MacDonald, A. H. Moiré bands in twisted double-layer graphene. Proc. Natl. Acad. Sci. USA 108, 12233–12237 (2011).
2. Cao, Y. et al. Correlated insulator behaviour at half-filling in magic-angle graphene superlattices. Nature 556, 80–84 (2018).
3. Po, H. C., Zou, L., Vishwanath, A. & Senthil, T. Origin of Mott insulating behavior and superconductivity in twisted bilayer graphene. Phys. Rev. X 8, 031089 (2018).
4. Lu, X. et al. Superconductors, orbital magnets and correlated states in magic-angle bilayer graphene. Nature 574, 653–657 (2019).
5. Zondiner, U. et al. Cascade of phase transitions and Dirac revivals in magic-angle graphene. Nature 582, 203–208 (2020).
6. Wong, D. et al. Cascade of electronic transitions in magic-angle twisted bilayer graphene. Nature 582, 198–202 (2020).
7. Cao, Y. et al. Unconventional superconductivity in magic-angle graphene superlattices. Nature 556, 43–50 (2018).
8. Yankowitz, M. et al. Tuning superconductivity in twisted bilayer graphene. Science 363, 1059–1064 (2019).
9. Stepanov, P. et al. Unravelling the insulating and superconducting orders in magic-angle graphene. Nature 583, 375–378 (2020).
10. Saito, Y., Ge, J., Watanabe, K., Taniguchi, T. & Young, A. F. Independent superconductors and correlated insulators in twisted bilayer graphene. Nat. Phys. 16, 926–930 (2020).
11. Sharpe, A. L. et al. Emergent ferromagnetism near three-quarters filling in twisted bilayer graphene. Science 365, 605–608 (2019).
12. Cao, Y. et al. Strange metal in magic-angle graphene with near Planckian dissipation. Phys. Rev. Lett. 124, 076801 (2020).
13. Polshyn, H. et al. Large linear-in-temperature resistivity in twisted bilayer graphene. Nat. Phys. 15, 1011–1016 (2019).
14. Abrikosov, A. A. & Khalatnikov, I. M. The theory of a Fermi liquid (the properties of liquid He at low temperatures). Rep. Prog. Phys. 22, 329–367 (1959).
15. Proust, C. & Taillefier, L. The remarkable underlying ground states of cuprate superconductors. Annu. Rev. Condens. Matter Phys. 10, 409–429 (2019).
16. Greene, R. L., Mandal, P. R., Poniatiowski, N. R. & Sarkar, T. The strange metal state of the electron-doped cuprates. Annu. Rev. Condens. Matter Phys. 11, 213–229 (2020).
17. Grigera, S. A. et al. Magnetic field-tuned quantum criticality in the metallic ruthenate Sr$_2$RuO$_4$. Science 294, 329–332 (2001).
18. Shibauchi, T., Carrington, A. & Matouda, Y. A quantum critical point lying beneath the superconducting dome in iron pnictides. Annu. Rev. Condens. Matter Phys. 5, 113–135 (2014).
19. Löhneysen, H. V. et al. Non-Fermi-liquid behavior in a heavy-fermion alloy at a magnetic instability. Phys. Rev. Lett. 72, 3262–3265 (1994).
20. Ghiotto, A. et al. Quantum criticality in twisted transition metal dichalcogenides. Nature 597, 345–349 (2021).
21. Bruin, J. A. N., Sakai, H., Perry, R. S. & Mackenzie, A. P. Similarity of scattering rates in metals showing $T$-linear resistivity. Science 339, 804–807 (2013).
22. Trovarelli, O. et al. YbB$_6$S$_6$: Pronounced non-Fermi-liquid effects above a low-lying magnetic phase transition. Phys. Rev. Lett. 85, 626–629 (2000).
23. Lacciardello, S. et al. Electrical resistivity across a nematic quantum critical point. Nature 567, 213–217 (2019).
24. González, J. & Stauber, T. Marginal Fermi liquid in twisted bilayer graphene. Phys. Rev. Lett. 124, 186801 (2020).
25. Efetov, D. K. & Kim, P. Controlling electron-phonon interactions in graphene at ultrahigh carrier densities. Phys. Rev. Lett. 105, 256805 (2010).
26. Wu, F., Hwang, E. & Sarma, S. D. Phonon-induced giant linear-in-$T$ resistivity in magic angle twisted bilayer graphene: Ordinary strangeness and exotic superconductivity. Phys. Rev. B 99, 165112 (2019).
27. Liu, X. et al. Tuning electron correlation in magic-angle twisted bilayer graphene using Coulomb screening. Science 371, 1261–1265 (2021).
28. Saito, Y. et al. Isospin Pomeranchuk effect in twisted bilayer graphene. Nature 592, 220–224 (2021).
29. Rozen, A. et al. Entropic evidence for a Pomeranchuk effect in magic-angle graphene. Nature 592, 214–219 (2021).
30. Zaanan, J. Why the temperature is high. Nature 430, 512–513 (2004).
31. Legros, A. et al. Universal $T$-linear resistivity and Planckian dissipation in overdoped cuprates. Nat. Phys. 15, 142–147 (2019).
32. Dao, R. et al. Linear temperature dependence of resistivity and change in the Fermi surface at the pseudogap critical point of a high- $T$ superconductor. Nat. Phys. 5, 31–34 (2009).
33. Hayes, I. M. et al. Scaling between magnetic field and temperature in the high-
temperature superconductor BaFe$_2$(As$_{1-x}$P$_x$)$_2$. *Nat. Phys.* **12**, 916–919 (2016).
34. Giraldo-Gallo, P. et al. Scale-invariant magnetoresistance in a cuprate
superconductor. *Science* **361**, 479–481 (2019).
35. Liao, Z. M., Zhou, Y. B., Wu, H. C., Han, B. H. & Yu, D. P. Observation of
both classical and quantum magnetoresistance in bilayer graphene.
*Europhys. Lett.* **94**, 57004 (2011).
36. Kisslinger, F. et al. Linear magnetoresistance in mosaic-like bilayer graphene.
*Nat. Phys.* **11**, 650–653 (2015).

**Publisher’s note** Springer Nature remains neutral with regard to jurisdictional claims in
published maps and institutional affiliations.
© The Author(s), under exclusive licence to Springer Nature Limited 2022
Methods

Screening layer fabrication process. The devices presented in this study were produced following the ´cut-and-stack´ method: a thin hexagonal boron nitride (hBN) flake is picked up with a propylene carbonate (PC) film and then placed onto a 90 °C polydimethyl siloxane (PDMS) stamp. The hBN flake then allows a portion of a pre-cut monolayer graphene flake previously exfoliated mechanically on a Si++/SiO₂ (285 nm) surface to be picked up. Then, the remaining graphene sheet is rotated to a target angle usually around 1.1–1.15° and picked up by an hBN/graphene stack on PC (introduced previously). The resulting heterostructure is placed on top of another thin hBN flake (whose thickness is chosen by optical contrast and further confirmed with atomic force microscopy measurements). Finally, the last layer of the heterostructure, sitting at the very bottom, is composed of a graphene flake (typically a few layers of graphene thick) to create both a local back gate and a screening layer. The final stack is then placed onto a target Si++/SiO₂ (285 nm) wafer, where it is etched into a multiple Hall bar geometry using CHF₃/O₂ plasma and edge-coupled to Cr/Au (5/50 nm) metal contacts.

Transport measurements. Transport measurements were carried out in a dilution refrigerator (with base temperature 20–40 mK) and an ‘Ice Oxford’ VTI (with base temperature 1.6 K). We used a standard low-frequency lock-in technique using Stanford Research SR860 amplifiers with excitation frequency f = 13.131 Hz and Stanford Research SR560 pre-amplifiers. The back-gate voltage was controlled using Keithley 2400s voltage source meters. Voltage (d.c.) versus excitation current (d.c.) measurements were also performed using an SR560 low-noise d.c. voltage preamplifier in combination with a Keithley 2700 multimeter. Superconducting-type coaxial cables (2 m long, Lakeshore) connected the mixing chamber plate to the room-temperature plate. Each line was equipped with a ∏/4 filter (RS-239-191) at room temperature, a powder filter (Leiden Cryogenics) and a two-stage resistor–capacitor filter at the mixing chamber plate to avoid any unwanted heating of the charge carriers at low temperature. Our measurements showed no sign of saturation of the electronic temperature down to 40 mK. Investigations to lower temperatures will require advanced filtering to avoid heating of the charge carriers and advanced thermometry techniques to accurately evaluate the electronic temperature. The measurements were performed for various excitation currents I (obtained by applying a tension to a 10 MΩ resistor), ranging from I < 10 nA, so as not to break the superconducting state, up to I = 200 nA (to enhance the signal in the metallic phase). Special attention was given to ensure consistency between datasets and not to increase the electronic temperature by gradually increasing the current and tracking the output voltage. All in-field measurements reported in this work were performed with an applied out-of-plane magnetic field.

Data availability

Source data are provided with this paper. All other datasets that support the plots within this publication are available from the corresponding authors upon reasonable request.

Acknowledgements

We are grateful for fruitful discussions with A. MacDonald, P. Jarillo-Herrero and P. Coleman. D.K.E. acknowledges support from the Ministry of Economy and Competitiveness of Spain through the ‘Severo Ochoa’ programme for Centres of Excellence in R&D (SEV-0522), Fundació Privada Cellex, Fundació Privada Mir-Puig, the Generalitat de Catalunya through the CERCA programme and funding from the European Research Council (ERC) under the European Union’s Horizon 2020 research and innovation programme (grant agreement number 852927). J.D.-M. acknowledges support from the INPhINIT ‘La Caixa’ Foundation (ID 100010434) fellowship programme (LCF/BQ/ID19/11730021). G.D.B. acknowledges funding from the ‘Presidencia de la Agencia Estatal de Investigación’ within the ‘Convocatoria de tramitación anticipada, correspondiente al año 2019, de las ayudas para contratos predoctorales (Ref. PRE2019-088487) para la formación de doctores contemplada en el Subprograma Estatal de Formación del Programa Estatal de Promoción del Talento y su Empleabilidad en IV+i+i, en el marco del Plan Estatal de Investigación Científica y Técnica y de Innovación 2017-2020, cofinanciado por el Fondo Social Europeo’. I.D. acknowledges support from the INPhINIT ‘La Caixa’ (ID 100010434) fellowship programme (LCF/BQ/ID19/11730030). K.W. and T.T. acknowledge support from the Elemental Strategy Initiative conducted by the MEXT, Japan (grant number JPMXP0112101001) and JSPS KAKENHI (grant numbers 19H05790 and JP20H00354). L.L. acknowledges support from the Science and Technology Center for Integrated Quantum Materials, NSF Grant No. DMR-1231319 and Army Research Office Grant W911NF-18-1-0116.

Author contributions

D.K.E. and X.L. conceived and designed the experiments. I.D., G.D.B., J.D.-M. and X.L. fabricated the devices. A.J., I.D., G.D.B., J.D.-M. and X.L. performed the measurements. A.J. analysed the data. A.J., I.D. and L.L. performed the theoretical modelling. T.T. and K.W. contributed materials. D.K.E. supported the experiments. A.J. and D.K.E. wrote the paper.

Competing interests

The authors declare no competing interests.

Additional information

Supplementary information The online version contains supplementary material available at https://doi.org/10.1038/s41567-022-01556-5.

Correspondence and requests for materials should be addressed to Alexandre Jaoui or Dmitri K. Efetov.

Peer review information Nature Physics thanks the anonymous reviewers for their contribution to the peer review of this work.

Reprints and permissions information is available at www.nature.com/reprints.