One-dimensional Luttinger liquids in a two-dimensional moiré lattice

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The Luttinger liquid (LL) model of one-dimensional (1D) electronic systems provides a powerful tool for understanding strongly correlated physics, including phenomena such as spin–charge separation1. Substantial theoretical efforts have attempted to extend the LL phenomenology to two dimensions, especially in models of closely packed arrays of 1D quantum wires2–11, each being described as a LL. Such coupled-wire models have been successfully used to construct two-dimensional (2D) anisotropic non-Fermi liquids2–6, quantum Hall states7–9, topological phases10,11 and quantum spin liquids12,13. However, an experimental demonstration of high-quality arrays of 1D LLs suitable for realizing these models remains absent. Here we report the experimental realization of 2D arrays of ID LLs with crystalline quality in a moiré superlattice made of twisted bilayer tungsten ditelluride (tWTe2). Originating from the anisotropic lattice of the monolayer, the moiré pattern of tWTe2 hosts identical, parallel 1D electronic channels, separated by a fixed nanoscale distance, which is tunable by the interlayer twist angle. At a twist angle of approximately 5 degrees, we find that hole-doped tWTe2 exhibits exceptionally large transport anisotropy with a resistance ratio of around 1,000 between two orthogonal in-plane directions. The across-wire conductance exhibits power-law scaling behaviours, consistent with the formation of a 2D anisotropic phase that resembles an array of LLs. Our results open the door for realizing a variety of correlated and topological quantum phases based on coupled-wire models and LL physics.

In various coupled-wire models2–13, one-dimensional (ID) quantum wires are placed in parallel with each other at an exactly fixed nanoscale distance, producing a two-dimensional (2D) or three-dimensional periodic system. In 2D, such perfectly arranged wires can, in principle, realize a strongly anisotropic non-Fermi liquid phase that resembles a Luttinger liquid (LL)2–5. When a perpendicular magnetic field is applied, new quantum Hall states7–9 may also develop in such an array without the presence of a free 2D electron gas. This highly anisotropic setting is qualitatively different from conventional isotropic 2D electron systems. Experimentally realizing these interesting coupled-wire constructions is challenging, as they require a large number of identical nanowires to be strictly arranged in a crystalline array at both nano and microscopic scales. A route to overcome these difficulties is to use moiré superlattices of a twisted bilayer stack of an anisotropic 2D crystal. Indeed, it has been proposed that twisted 2D crystals with a rectangular unit cell, such as GeSe14, create 1D flat bands. Another excellent choice is tWTe2, as its monolayer unit cell is an elongated rectangle. In this work, we uncover the potential of tWTe2 for creating the desired high-quality arrays of 1D wires that can expand the LL physics to 2D.

tWTe2, moiré lattices and device design

Monolayer WTe2 consists of three atomic layers (Te-W-Te) in a sandwich structure, in which the W atoms are organized in ID zigzag chains15 (Fig. 1a). The tWTe2 has six atomic layers with a complicated moiré pattern. To better illustrate the moiré lattice of small-angle tWTe2, we present the superlattice of only the W layers and of only the Te layers separately in Fig. 1b, c. The Te pattern develops a triangular superlattice viewed from the top, whereas the W moiré pattern develops ID stripes, reflecting the underlying anisotropy of the monolayer. The bright stripes in Fig. 1b indicate regions where the W atoms from two layers are optimally aligned vertically (AA stripes), whereas in the dark stripes they are optimally misaligned (AB stripes). The distance d between neighbouring AA stripes depends on the twisted angle θ, d = a/2sin(θ/2)), for small θ (Fig. 1d); here a is the length of the monolayer’s rectangular unit cell. In Fig. 1e and Extended Data Fig. 1, we experimentally visualize this unique moiré structure of tWTe2 using conductive atomic force microscopy (cAFM). Below we present transport studies of two devices with θ = 5° (d = 7.2 nm, device no. 1) and θ = 6° (d = 6.0 nm, device no. 2).

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Figure 1 illustrates the design of our devices for measuring the transport properties of tWTe₂. Similar to our previous reports, a thin hexagonal boron nitride (hBN) layer is inserted between the tWTe₂ and the palladium (Pd) electrodes, with selected areas etched in the thin hBN layer to expose the very ends of Pd for contacting the tWTe₂. Such a device geometry restricts the contact area to be small in the 2D bulk and eliminates transport contributions of conducting bulk. To avoid electrical contact to non-tWTe₂ regions, a hBN layer with selectively etched areas is used.

**Exceptionally large transport anisotropy**

We first examine transport anisotropy in the tWTe₂ devices. Figure 1i shows plots of the four-probe resistance $R_{tg}$ as a function of gate-induced doping $n_g$, taken from device no. 1 at 1.8 K with the contact configurations $(R_{hard} and R_{easy})$ shown in Fig. 1h. Here $n_g = \epsilon, s, d (V_{tg}/d_{bg} + V_{bg}/d_{tg})/\epsilon$, where $d_{bg}$ is the thickness of hBN dielectric layers for the top (bottom) gate; $V_{tg}(V_{bg})$ is the top (bottom) gate voltage; $\epsilon, s, d$ are the elementary charge, vacuum permittivity and relative dielectric constant of hBN, respectively. The current is applied along two orthogonal directions in the atomic plane for measuring $R_{hard}$ and $R_{easy}$, respectively. The choice of the easy and hard directions was made by comprehensively examining two-probe resistances taken between all neighbouring electrodes and other configurations along and across moiré stripes (Extended Data...
Conductance power laws

The strongly anisotropic phase of tWTe\textsubscript{2} exhibits robust power-law and scaling behaviours in the across-wire transport (Fig. 2b), in which currents flow perpendicular to the moiré stripes (that is, the wires). Figure 2c plots the measured across-wire conductance \(G\) at a selected gate voltage in the hole-doped regime, with a two-probe configuration shown in Fig. 2b. As seen in the log–log plot, \(G \propto T^\alpha\), for \(T\) below about 30 K, with an exponent \(\alpha\) of approximately 0.98 (or approximately 1.14) for device no. 1 (or no. 2) at the chosen gate voltages. To demonstrate scaling, we present differential conductance \((dG/dV)\) measurements under varying both the d.c. source-drain bias \((V)\) and \(T\) (Fig. 2d). For small enough bias, \(dG/dV\) develops plateaux, indicating that the conductance is controlled only by \(T\) via the power law. At high bias, all curves taken at different \(T\) merge together, with a trend that can be well captured by the 'same' power-law exponent \(\alpha\), that is, \(dG/dV \times V^\alpha\). Indeed, in the scaled conductance plot (Fig. 2e), \((dG/dV)/T^\alpha\) versus \(eV/\hbar\)\(g\)\(T\), all data points, taken in a parameter range wider than a decade in \(T\) and three decades in \(V\), collapse into a single curve. A similar collapse can be found in device no. 2 (Fig. 2f) and Extended Data Fig. 7.

In Extended Data Fig. 8, we compare along-wire and across-wire conductance taken from device no. 1. More robust power laws are typically seen in the across-wire direction. We note that contact resistance plays an important role in the along-wire transport, as seen in nanotubes. In practice, the moiré system in the contact region may be affected by distortions, strain, unintentional doping and other interface effects. By contrast, across-wire resistance at low \(T\) is dominated by the tWTe\textsubscript{2} bulk (see Extended Data Fig. 9), a much more uniform area.
Gate-tuned anisotropy cross-over

The power-law across-wire conductance is generally observed over a wide range of gate voltages for $T < 30$ K in the tTe samples, as illustrated in Fig. 3 and Extended Data Figs. 10–12. We extract the $n_g$ dependence together with $\beta_g$ for devices no. 1 and no. 2, respectively (Fig. 3a, e). In the hole-doped side, strong anisotropy occurs together with good power-law scaling characteristics, as shown by the collapse of the $dI/dV$ curves over a wide range of $T$ and $V$ in the scaled plots (Fig. 3b–d for device no. 1; Fig. 3f–h for device no. 2). Although in both devices $\alpha$ is valued near unity on the hole side, the exact gate-dependent behaviour differs, which could arise owing to twist-angle-dependent electronic structures or extrinsic effects, such as disorders. Near charge neutrality or on the electron-doped regime, high-bias data deviate from the power-law trend (Extended Data Figs. 11 and 12). With electron doping, transport anisotropy is strongly suppressed, although zero-bias $G(T)$ still approximately follows a power law with a decreasing exponent down to near zero at high electron doping (Extended Data Fig. 10).

Band structure modelling

We further perform a continuum model analysis on tTe at the single-particle level (Fig. 4). The modelling is challenging, as even at the monolayer level, topology20–23, correlations16,24–25 and spin–orbit coupling20 are all present. We start with a density functional theory (DFT) calculation on the monolayer, yielding valence band maximum at $\Gamma$ flanked by two conduction band minima at the wavevector $\pm q$. (Fig. 4a). DFT calculations for untwisted but shifted bilayers are used to extract effective interlayer couplings (Fig. 4b, c), which enter the continuum model for obtaining the tTe structure26. The resulting twisted bands arising from one conduction valley are shown in Fig. 4d, in which a pair of highly anisotropic bands indeed develop in the hole regime, in contrast to the electron regime for which no substantial anisotropy is seen. Figure 4f illustrates the corresponding quasi-1D hole Fermi surface with the corresponding real space wavefunctions that coincide with the moiré stripes (Fig. 4e). In contrast to the $\pm q$ valleys, the moiré reconstruction of the valence bands at $\Gamma$ is much less pronounced in this simplified model and develops no large anisotropy (Methods). Note that, although our simple analysis here does capture the emergence of quasi-1D bands, a comprehensive modelling would necessarily require future efforts involving large-scale DFT calculations, lattice reconstructions and interaction effects.

The Luttinger liquid interpretation

The large resistance anisotropy and contrasting $T$ dependence of $R_{\text{hard}}$ and $R_{\text{easy}}$ (Extended Data Fig. 4) indicate that transport is qualitatively different between across- (insulating) and along-wire (metallic) directions. The power-law behaviour itself is inconsistent with the formation of an ordinary band or Mott insulator. The exponent $\alpha$ varies smoothly as a function of $n_g$, showing no obvious presence of a fully insulating state (Fig. 3), consistent with the absence of a gap in the modelling. For a 2D diffusive metal27,28, the ‘tunnelling anomaly’ owing to the relaxation of injected charges at the contact may lead to a conductance power law depending on $T$ or $V$. However, this cannot give distinct transport exponents in different directions and cannot account for our observation along the hard direction, where the resistance is dominated by the tTe bulk rather than contact effects. For disordered quasi-1D systems, calculations have shown that variable...
range hopping transport may produce an apparent power-law behaviour, that is, \( G \propto T^a \) for \( eV \ll k_B T \) and \( \frac{d}{dt}V \propto V^{\beta} \) for \( k_B T \ll eV \), where \( \alpha \) and \( \beta \) are two generally unequal exponents that are independently controlled by microscopic details including disorder. This is, however, in sharp contrast to our observation in tWTe\(_2\), where the power laws in \( T \) and \( V \) are controlled by the same exponent, namely, \( \alpha = \beta \). This single-exponent scaling behaviour is robustly observed over a wide hole-doping range (Fig. 3 and Extended Data Figs. 11 and 12), where \( \alpha \) has been tuned, and in samples with varied twist angles (devices no. 1 and no. 2). These observations provide strong evidence that the single-exponent power-law behavior is generic to the anisotropic phase of tWTe\(_2\). Any explanation that requires fine-tuning of parameters to achieve the condition of \( \alpha = \beta \) is unlikely to be feasible.

A natural explanation is emergent LL physics. The characteristic feature of a 1D LL in transport is indeed the 'single-exponent' power-law dependence of its conductance, that is, \( \alpha = \beta \). The power-law transport of LL physics has been observed in several 1D systems, such as nanotubes\(^{40,43,41}\), engineered 1D structures\(^{31-33}\), edge modes\(^{34-36}\), polymers\(^{37}\) and self-organized gold wires\(^{38}\). However, extending LL physics from a single 1D wire to a 2D system is of fundamental interest yet challenging. Proposals to do so on the basis of 2D arrays of 1D wires have been put forward\(^{2,4}\), but, as far as we know, the intriguing concept of a 2D anisotropic phase that mimics a LL has so far not been established in real materials.

Our observations on the hole side of tWTe\(_2\) are well consistent with the generic LL expectations. We therefore propose that the anisotropic phase arises because of the formation of a 2D array of 1D LLs induced by the moiré superlattice. Understanding the moiré-induced LL behaviours in tWTe\(_2\) requires proper consideration of electron interactions and transport mechanisms. In a quasi-1D system, although the early calculations\(^{39}\) of the across-wire transport exponent between parallel LLs indicated \( \alpha = 2\eta \), where \( \eta \) (which vanishes without interactions) is the Fermi surface exponent for an individual wire determined by the LL parameter \( K \), the more recent consensus is instead that \( \alpha = 2\eta - 1 \), where the extra −1 arises from the fact that hopping can occur anywhere along the wires\(^{40}\). The relationship applies when single-particle hopping is the dominant conduction process and \( T \) is much larger than the 1D to 2D cross-over temperature \( T^* = \frac{t_\parallel}{t_\perp}/(1 - \frac{\eta}{2}) \approx 50 \) meV, respectively. If this applies, for certain hole doping of tWTe\(_2\), the across-wire conductance exhibits a power-law exponent \( \alpha = 1 \), corresponding to an effective \( \eta = 1 \), near the marginal boundary above which the single-particle process is no longer relevant and two-particle processes may be important. Assuming spin degeneracy and \( \eta = (K + 1/K - 2)/4 \) (ref. 40), we obtain an effective intrawire \( K = 0.17 \) for \( \alpha = 1 \). The strong intrawire interaction is consistent with the experimental fact that the deviation from the power law is absent down to at least 1.8 K. This remarkably stable LL behaviour in the anisotropic 2D system calls for careful consideration of the interaction-driven phases in tWTe\(_2\), especially the interwire interactions given the nanoscale wire spacing. Considering interwire interactions, the transport exponent then depends on a stiffness function \( x(q_z) \) instead of a single intrawire parameter \( K \), where \( q_z \) is the momentum perpendicular to the wires\(^{2,4}\). We note that further experimental and theoretical explorations are necessary to examine the exact connection between the measured power laws to interactions in the system, which is critical to evaluate the enticing possibility of a sliding LL phase and a host of competing orders descending from it\(^{2,4}\).
Summary

We demonstrate a new tuneable platform based on tWTe₂ stacks for studying high-quality 2D arrays of 1D electronic structures in a crystalline superlattice. We interpret the results on the basis of the formation of a 2D anisotropic non-Fermi liquid phase that resembles a LL. An exciting direction is to search for new quantum Hall states with an applied magnetic field. The physics of spin–charge separation, naturally expected in LLs, is another interesting direction to pursue. Experimental searches for evidence of spin–charge separation in a 2D WTe₂ system could provide important opportunities for studying new regimes in strongly correlated quantum phases.

Online content

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**Methods**

**Sample fabrication**

We followed the WTe$_2$ crystal growth, exfoliation and device fabrication procedures detailed in refs. 15–17,24. The stack of WTe$_2$ was obtained by the ‘tear and stack’ technique$^{42,43}$, in which we used the hBN layer to tear and pick up part of a monolayer WTe$_2$ flake, followed by rotating the rest of the flake counterclockwise by a chosen angle $\theta$ and then stacking the two monolayer WTe$_2$ pieces together. We summarize a step-by-step fabrication process in Extended Data Fig. 2. The device for CAFM measurements consists of a stack of few-layer hBN/tWTe$_2$/hBN on top of a layer of Pd thin film.

**Conductive atomic force microscopy measurement**

The CAFM measurements$^{44}$ were performed at room temperature in a Bruker Dimension Icon AFM with dry nitrogen purged into the acoustic shield to eliminate the oxygen and water and reduce WTe$_2$ degradation. A humidity as low as <0.1% and no more than 5% was kept during the measurements. A PF-TUNA module equipped with an in situ current amplifier was used. The device was biased with $\pm 2.5$ mV d.c. voltage and the d.c. current through the tip was recorded. The 2D conductance image was then captured and plotted, as shown in Extended Data Fig. 1b, c. To better view the moiré pattern, we have applied a standard flatten process and filtered out the 60 Hz electronic noise to the AFM image.

**Transport measurement**

The electrical measurements of our devices were performed in a cryostat (Quantum Design Dynacool) equipped with a superconducting magnet. Standard lock-in measurements were taken with a frequency of 3–78 Hz. The four-probe measurements were performed by supplying an a.c. current of approximately 3 nA. The two-probe differential conductance measurements were carried out with a small a.c. excitation of 50 or 100 µV, together with a d.c. excitation up to about 400 mV. A current pre-amplifier (DL Instrument 121I) and a voltage pre-amplifier (DL Instrument 120I) were used to improve signals. Two Keithley 2450 source meter units were used to control the top and bottom gates.

**Estimating intrinsic anisotropy from four-probe measurements**

Following ref. 45, we discuss the impact that sample geometry and contact placement have on four-probe resistance measurements. In particular, we are interested in how an intrinsic sheet resistivity anisotropy $\rho_{\text{bulk}} = \rho_{x\gamma} = \rho_{y\gamma}$ translates to an observed four-probe anisotropy $\rho_{\text{top}} = R_{x\gamma} / R_{y\gamma}$. We address this by considering the electric potential distribution over the sample in the classical limit. To simplify the problem, we assume that: (1) no current leaks out of the tWTe$_2$ sample boundary; (2) the sample is characterized by a spatially uniform resistivity tensor (2) the external current source/sink distribution $f(x,y)$ is modelled as delta functions at the current contacts; and (4) the voltage (and unused) contacts do not substantially change the physics. By combining the continuity equation, Ohm’s law and Faraday’s law, we obtain an anisotropic Poisson equation for the scalar potential with derivative boundary conditions

$$\left[\frac{1}{\rho_{\text{bulk}}} \partial_x^2 \phi(x,y) + \frac{1}{\rho_{\text{bulk}}} \partial_y^2 \phi(x,y)\right] = -\rho_{\gamma\gamma} f(x,y)$$

$$\frac{1}{\rho_{\text{bulk}}} \partial_x \phi(x,y), \partial_y \phi(x,y) \right) \cdot \vec{n}(x,y) = 0 \quad \text{for} (x,y) \in \partial S$$

where $S$ denotes the WTe$_2$ region. From this, the measured resistance can be obtained once the voltage contact locations are prescribed.

We consider a simple caricature of the experimental set-up by taking the sample to be a square of side length $L$, aligned with the principal axes of the resistivity tensor, that is, the region $[-L/2,L/2] \times [-L/2,L/2]$. For a measurement of $R_{\gamma\gamma}$, the current contacts are placed at $(\pm L/4,0)$ and the voltage contacts at $(\pm L/8, L/4)$. $R_{\gamma\gamma}$ is computed similarly, except the contacts are rotated by $90^\circ$. The results are shown in Extended Data Fig. 5 for different values of the $\rho_{\text{bulk}}$. Above some critical value, the measured anisotropy $\rho_{\gamma\gamma}$ grows exponentially as the square root of $\rho_{\text{bulk}}$.

**Scaling formula and the fitting procedure**

To fit the conductance data shown in Fig. 2, we derive a scaling formula for across-wire transport by following the procedure and assumptions in ref. 46. The resulting formula is:

$$\frac{1}{T^2} \frac{dI}{dV} = A \left[ \cosh \left( \frac{1 + \alpha}{2} \frac{eV}{kT} \right) + \frac{1}{\pi^2} \left( \frac{eV}{kT} \right)^2 \right] \Gamma(z) + \frac{\sinh \left( \frac{1 + \alpha}{2} \frac{eV}{kT} \right)}{\gamma} \left( \frac{eV}{kT} \right)^2 \left( \frac{1}{\pi^2} \Gamma(z)^2 + \left( \psi(z) + h.c. \right) \right)$$

where the dimensionless variable $x = \frac{eV}{kT}, \alpha = \frac{1 + \alpha}{2} \frac{eV}{kT}$, $z = \frac{1 + \alpha}{2} \frac{eV}{kT}$, $\Gamma(z)$ is the gamma function, $\psi(z)$ is the digamma function, $\alpha$ is the power-law exponent, $\gamma$ is a constant introduced to account for the division of the source-drain voltage across the multiple wires in series, $A$ is an overall coefficient and $\gamma$ is the Hermitian conjugate. This formula assumes that the dominant transport mechanism between parallel LL channels is single-particle hopping$^{46}$. Note that we have taken a derivative of $I(V)$ to obtain an expression for $dI/dV$. In the fitting procedure, we first assign the $a$ and $1/y$ to specific values and find the best fit by optimizing the parameter $A$; for each combination of $a$ and $1/y$, a root mean squared error (r.m.s.e.) considering all data points is calculated. The r.m.s.e. as a function of $a$ and $1/y$ is plotted as the 2D colour plot (see details in Extended Data Fig. 7 for both device no. 1 and no. 2). The best fits fall in the regions with minimized r.m.s.e. in the 2D plots. We note that our modelling of the WTe$_2$ system is at a very early stage at present, and hence new scaling formulae that better describe WTe$_2$ transport may be developed in the future, with improved understanding of the system. However, we emphasize that the key experimental demonstration of the power-law scaling here is based on the direct observation of the single-exponent behaviour, that is, $G \propto \gamma^\alpha$ at the low bias limit, $dI/dV \propto \gamma^\alpha$ at the high-bias limit (Fig. 2d) and the fact that all data collapse in the scaled conductance plot (Fig. 2e). These key features are independent of any fitting formula used for analysis.

**DFT calculations**

DFT calculations on untwisted systems were performed using the plane-wave pseudopotential code QUANTUM ESPRESSO$^{47}$. For the band structure of monolayer WTe$_2$ we used fully relativistic optimized norm-conserving Vanderbilt pseudopotentials from PseudoDojo$^{48}$ and the PBE exchange-correlation functional$^{49}$ with a $\times 6 \times 1 \times 1$ k-grid. The van der Waals corrections were included via the semi-empirical framework of DFT-D3 (ref. 50). The plane-wave cut-off was 80 Ry and Marzari–Vanderbilt smearing of width 0.01 Ry was used. We consider a cell of height 35 Å to eliminate the effect of periodic images. The fully relaxed monolayer geometry was obtained without spin–orbit coupling (SOC) and van der Waals corrections, with a force tolerance of 10$^{-9}$ atomic units on each atom. The shifted untwisted bilayers were constructed by fixing the in-plane positions of the atoms in one of the layers to be displaced by $d = (dx,dy, dz)$, but letting the out-of-plane coordinates freely relax. The resulting band structures, obtained without SOC, determine the effective interlayer couplings used in the continuum models$^{34}$. The neglect of SOC in the bilayers is a good approximation, as the band splitting is mostly due to interlayer interactions and is relatively insensitive to SOC. Furthermore, the bands remain spin degenerate, as inversion symmetry is retained for any $d$. Although lattice relaxation is not considered, we expect its effect would be important in real devices as it will
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deform the heterostructure by expanding the low-energy stackings at the expense of the high-energy regions.

Continuum model calculations

Moiré continuum models can be constructed using input from monolayer and untwisted bilayer DFT data. Doing this for twisted WTe₂ is challenging because (1) the unit cell is rectangular, (2) the (multiple) bands near the Fermi energy are complicated and (3) some of the relevant low-energy features generally lie away from high-symmetry points. This is to be contrasted with hexagonal-based systems, such as graphene (Dirac cones at K), or other transitional metal dichalcogenides (parabolic bands with large band gaps at Γ or K). To make progress, we take the small-angle limit and assume that the main features of monolayer WTe₂ can be treated separately (that is, the valence band maximum and conduction band minima (valleys)). Although these bands have significant energetic overlap, for small twist angles and smooth moiré potentials the coupling between these bands will be suppressed.

The general approach is to first model an effective Bloch Hamiltonian \( H(k, d) \) that describes the untwisted bilayer for different in-plane shifts \( d \) (it is assumed here that the out-of-plane coordinates have been relaxed). This typically involves kinetic energy terms \( E(k) \) from the individual layers, as well as (layer-dependent) interlayer potential \( V(k, d) \) and hopping terms \( \Delta(k, d) \). If we take one band from each layer, the effective Hamiltonian is

\[
H(k, d) = \left( \begin{array}{cc} E(k) + V_0(k, d) & \Delta(k, d) \\ \Delta^*(k, d) & E(k) + V_0^*(k, d) \end{array} \right)
\]

where \( \Delta^*(k, d) \) is the complex conjugate of \( \Delta(k, d) \). Then the relation \( d = \theta x r \), valid for rigid twist, is used to convert the interlayer interactions into ‘hopping’ terms in the moiré reciprocal lattice. Note that, depending on the momentum base point, the twist will also induce a relative difference between the kinetic terms in the two layers. In particular, the conduction valley continuum model is based at \( q_v \), which is around one-third the distance to the Brillouin zone boundary in the \( \alpha \) direction. The resulting matrix (in combined band/layer/moiré reciprocal lattice vector space) is diagonalized with a large enough plane-wave cut-off to ensure convergence in the energy window of interest.

We apply the approach above to obtain the reconstructed moiré bands from a conduction band valley. SOC is neglected (as in the bilayer DFT calculations) and spin degeneracy is assumed (inversion is only valid for rigid twists, is used to convert the interlayer interactions into ‘hopping’ terms in the moiré reciprocal lattice).

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Additional information

Supplementary information

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Extended Data Fig. 1 | Conductive atomic force microscope (cAFM) measurements of tWTe₂ at room temperature. 

**a**, Cartoon illustrations of the cAFM measurement on a few-layer hBN/tWTe₂/hBN stack on a Pd metal film pre-patterned on a SiO₂/Si chip. The inset illustrates the cross-section of the stack. A relatively thick (around 39 nm) hBN was used to mitigate the roughness of the metal surface. 

**b**, A cAFM image taken from a θ ≈ 5° tWTe₂ device, directly visualizing the moiré structure. The dashed-dot square locates a zoom-in scan, as shown in **c**. We comment on three aspects of the observations. (1) The measurement was taken at room temperature, where the transport shows no significant anisotropy. This is consistent with the fact the measured local conductance \( G \) varies only by a small amount at different tip locations in the map. (2) As shown in **b** and **c**, the small variations already allow us to clearly image the underlying moiré structure. (3) Our experimental resolution does not allow us to identify which one is the AA stripe or the AB stripe, but the map is clearly consistent with the pattern shown in Fig 1b, except with lattice relaxations. The relatively low and high conductance regions develop into stripes, with an inter-stripe distance approximately 7.1 nm, consistent with the expectation for a θ = 5° tWTe₂ stack. (4) As WTe₂ is air sensitive, we have to use a few-layer hBN as a protecting layer and a sample fabrication process that minimizes the time for the sample exposed to air. The top surface of the thin hBN is left behind with polymer residues etc., which we believe could be the main source of the residue-like features in **b** and **c**. Our transport devices (devices no. 1 & no. 2 in the main text) use a top graphite gate that serves as a screening layer and hence the tWTe₂ channel is of much higher quality. Other details about the cAFM measurements can be found in the Methods.
Extended Data Fig. 2 | Sample fabrication process. **a**, Cartoon illustrations of a prepared top graphite/hBN stack (top) and a flake of monolayer WTe₂ on a SiO₂/Si chip (bottom). Their corresponding optical images are also shown to the right. The distance between the aligned graphite edge and hBN edge is carefully optimized during transfer to be within 500 nm. **b & c**, Tear the monolayer WTe₂ by the hBN into two separate pieces, labelled as top and bottom WTe₂. **d**, Rotate the bottom WTe₂ flake, i.e., the SiO₂/Si chip, counterclockwise by θ. **e**, Pick up the bottom WTe₂ to create a tWTe₂ stack. The optical image of the tWTe₂ stack shown to the right was taken after flipping the stamp upside down. The tWTe₂ stack is highlighted by the red dashed line. **f**, The pre-patterned bottom part (thin hBN/Pd electrodes/hBN dielectric/graphite) with etched holes in the thin hBN layer to expose the tips of the Pd electrodes (see ref. 16,17). An atomic force microscope image shows a clean surface of a prepared bottom stack. **g**, The final stack of a tWTe₂ device, with an optical image of a final device (no. 1). The thickness of flakes is 7.6 nm graphite/8.9 nm top hBN/WTe₂/2.0 nm thin hBN/5.5 nm bottom hBN/7.0 nm graphite for device no. 1 and 4.8 nm graphite/10.6 nm top hBN/WTe₂/4.2 nm thin hBN/15.0 nm bottom hBN/6.9 nm graphite for device no. 2. All scale bars are 3 μm.
Extended Data Fig. 3 | More analysis on the transport anisotropy on device no. 1.  

a, Two-probe resistance between neighbouring electrodes as a function of \( n_g \) in cooldown no. 1. Inset shows the contact configurations for each measurement, where the estimated hard direction (stripe direction) is indicated by the grey lines (not to scale). \( R_{2-3} \) and \( R_{6-7} \) display larger values than all others in the hole-doped regime, signifying the hard direction, whereas \( R_{4-5} \) shows the lowest value. Contact 1 was broken during the fabrication. The contact resistance plays a significant role here. After the easy and hard directions were identified, we performed four-probe measurements, as shown in Fig. 1 in the main text.

b, Two-probe resistance across (\( R_{2-3} \), \( R_{6-7} \), and \( R_{5-8} \)) and along (\( R_{4-5} \), \( R_{2-7} \), and \( R_{3-6} \)) the stripes as a function of \( n_g \) in cooldown no. 2. Inset shows the contact configurations for each measurement, where the easy direction (along stripes) is indicated by the grey lines (not to scale).

c, Four-probe resistance across (\( I_{sd}: 5-8; V_{xx}: 6-7 \) and \( I_{sd}: 5-8; V_{xx}: 3-2 \)) and along (\( I_{sd}: 3-6; V_{xx}: 4-5 \), \( I_{sd}: 3-6; V_{xx}: 2-7 \), and \( I_{sd}: 2-7; V_{xx}: 4-5 \)) the stripes as a function of \( n_g \) in cooldown no. 2.
Extended Data Fig. 4 | n_g dependent transport anisotropy for devices no. 1 and no. 2. Four-probe resistance R_xx as a function of n_g measured with an excitation current applied along hard and easy directions in linear plots. a, data taken for device no. 1 at 1.8 K (the same data as Fig. 1i). b-d, R_hard, R_easy, and R_hard/R_easy as a function of n_g at different temperatures taken from device no. 1. e-h, The same plots for device no. 2.
Extended Data Fig. 5 | Electrostatic simulation for the four-probe contact configuration. a & b, Electric potential distribution for contact arrangements corresponding to $R_{\text{hard}}$ and $R_{\text{easy}}$ four-probe measurements respectively (see Methods). Black dots indicate current contacts that source/sink current. Red dots indicate the placement of voltage contacts. c, Predicted four-probe anisotropy $\beta_{4p} = R_{\text{hard}} / R_{\text{easy}}$ as a function of the intrinsic sheet resistivity anisotropy $\beta_{\text{bulk}}$. For $\beta_{4p} = 1000$, we estimate $\beta_{\text{bulk}} \approx 50$. 
Extended Data Fig. 6 | Dual-gate-dependent transport along hard and easy directions (device no. 1, cooldown no. 1). The four-probe resistance taken at 1.8 K (200 K) along the hard and easy directions were shown in a (d) and b (e), respectively. $R_{\text{hard}}/R_{\text{easy}}$ at 1.8 K (200 K) is shown in c (f).
Extended Data Fig. 7 | Fitting to the differential conductance data based on the universal scaling formula. a, Raw data for the differential conductance measurements taken in device no. 1 (replotted from Fig. 2d). b, 2D map of calculated root-mean-square error (r.m.s.e.) as a function of the fitting parameters, $\alpha$ and $1/\gamma$ (see Methods for details). The best fit is obtained by finding the minimal value of r.m.s.e. in this plot, i.e., $\alpha = 0.94$ and $1/\gamma = 9$. c & d, Scaled conductance as a function of scaled excitation by assigning $\alpha = 0.94$ in a log-log plot and log-linear plot. The dashed line indicates the fitting result given by the universal formula defined in the Method section. e-h, The same fitting plots for device no. 2, using the same raw data shown in Fig. 2f.
Extended Data Fig. 8 | Comparison between along-wire and across-wire transport. 

**a**, Illustration of tWTe₂ moiré stripes on the electrodes (top view). **b**, Illustration of transport along wires. At low T, the along-wire transport is dominated by contact resistance, i.e., tunnelling from the metal (FL) to the moiré wires (LL). **c**, Illustration of the across-wire transport, where the dominant resistance is due to interwire tunnelling in the stripe regime (i.e., LL to LL tunnelling). **d**, Along-wire two-probe conductance G as a function of T, plotted in log-log scale at a selected gate parameter. A power-law fit (solid line) to the low T data is shown. **e**, Differential conductance dI/dV taken under the along-wire transport configuration as a function of d.c. bias V at different T. The dashed line indicates a power-law trend. The dot-dash line indicates a deviation from the trend at high bias. Note that distortions, strain, unintentional doing and other interface effects occur at the moiré in the contact regime, which could cause the deviation. **f & g**, the same plot for data taken from the across-wire transport (the same data as Fig. 2c, d), exhibiting a more robust power-law behaviour to higher bias and T. This can be understood as the dominant resistance in the across-wire transport comes from the tWTe₂ channel regime, which is more uniform compared to the contact regime. Data were taken from device no. 1 in cooldown no. 1.
Extended Data Fig. 9 | Comparison of two-probe and four-probe measurements across the wires. Cartoon illustration of (a) two-probe ($G_{2p}$) and (b) four-probe ($G_{xx}$) configurations used for the measurements. (c) and (d), $G_{2p}$ and $G_{xx}$ as a function of temperature taken in the hole-doped region ($n_g = -5.5 \times 10^{12} \text{ cm}^{-2}$ and $n_g = -13.5 \times 10^{12} \text{ cm}^{-2}$, respectively). At low $T$ (1.8 K - 25 K) the trends of $G_{2p}$ and $G_{xx}$ both follow a power law and match well, demonstrating that the power law is intrinsic to the tWTe$_2$ channel. At high $T$, the two trends of $G_{2p}$ and $G_{xx}$ deviate from each other, which can be understood as $G_{2p}$ saturates due to contact resistances whereas $G_{xx}$ is strongly affected by the temperature induced changes of anisotropy. The effective geometry factor, important for determining $G_{xx}$, changes as the sample is tuned from a strongly anisotropic phase at low $T$ to an isotropic phase at high $T$. The main analyses in this paper are focused on the low $T$ regime. The measurements were performed on device no. 1 in cooldown no. 2.
Extended Data Fig. 10 | Gate-tuned anisotropy cross-over. a, The across-wire two-probe conductance $G(T)$ displays a power-law relation ($G \propto T^\alpha$) for a wide range of doping for device no. 1 (cooldown no. 2). The colour of the data points encodes $n_g$, as shown in the colour bar. The solid lines are the power-law fittings, where the extracted exponent $\alpha$ is shown in the inset. The grey line replots the anisotropy ratio. b, The same plots for device no. 2 (cooldown no. 1). The grey line replots the anisotropy ratio shown in the inset of Fig. 1h. c, The same plots for device no. 2 (cooldown no. 2). Note that data taken from two different cooldowns from device no. 2 shows qualitatively consistent results with only minor quantitative differences (dashed line in the inset of c is the exponent $\alpha$ replotted from the inset of b for comparison). Arrows to the insets in a and c indicate the selected $n_g$, at which the scaling analysis of the differential conductance is performed in Extended Data Figs. 11 and 12, respectively.
Extended Data Fig. 11 | Additional power-law scaling analysis for device no. 1 (cooldown no. 2). The corresponding \( n_g \) for each data set is indicated in the inset of Extended Data Fig. 10a.  

- **a**: Temperature dependent across-wire two-probe conductance \( G(T) \) taken at the indicated \( n_g \). The solid line is the power low fit. \( G \propto T^\alpha \); \( \alpha = 1.15 \).

- **b'**: Bias dependent differential conductance taken at the same \( n_g \) under different \( T \). The dashed line indicates the power-law trend with the same exponent \( \alpha \) extracted in a. \( \alpha = 1.06 \).

- **b''**: The same data in b', but replotted as scaled differential conductance \( (dI/dV)/T^\alpha \) v.s. scaled bias \( eV/k_B T \). Other panels are the same plots for different \( n_g \). As seen in the plots, in the hole side (a-c) the data generally follows a power law very well, whereas near charge neutrality (d) and in the electron side (e), deviations start to develop at high bias. In the highly electron-doped region (f), \( dI/dV \) and \( G \) vary only a little bit (\( \alpha \approx 0 \)) with changing both \( V \) and \( T \), hence the behaviour is approximately ohmic. Data used for Figs. 3b-d are indicated in the lowest panel.
Extended Data Fig. 12 | Additional power-law scaling analysis for device no. 2 (cooldown no. 2). The corresponding $n_g$ for each data set is indicated in the inset of Extended Data Fig. 10c. a, Temperature dependent across-wire two-probe conductance $G(T)$ taken at the indicated $n_g$. The solid line is the power law fit. $a'$, Bias dependent differential conductance taken at the same $n_g$ under different $T$. The dashed line indicates the power-law trend with the same exponent $\alpha$ extracted in $a$. $a''$, the same data in $a'$, but replotted as scaled differential conductance $(d/dV)/T$ v.s. scaled bias $eV/k_BT$. Other panels are the same plots for different $n_g$. As seen in the plots, in the hole side (a-d) the data generally follows a power law very well, whereas near charge neutrality (e), deviations start to develop at high bias. In the highly electron-doped region (f), $dI/dV$ and $G$ vary only a little bit ($\alpha \approx 0$) with changing both $V$ and $T$, hence the behaviour is approximately ohmic. Data used for Figs. 3f–h are indicated in the lowest panel.