Magic-angle twisted trilayer graphene (TTG) has recently emerged as a platform to engineer strongly correlated flat bands. We reveal the normal-state structural and electronic properties of TTG using low-temperature scanning tunneling microscopy at twist angles for which superconductivity has been observed. Real trilayer samples undergo a strong reconstruction of the moiré lattice, which locks layers into near-magic-angle, mirror symmetric domains comparable in size with the superconducting coherence length. This relaxation introduces an array of localized twist-angle faults, termed twistons and moiré solitons, whose electronic structure deviates strongly from the background regions, leading to a doping-dependent, spatially granular electronic landscape. The Fermi-level density of states is maximally uniform at dopings for which superconductivity has been observed in transport measurements.

The prediction of a magic angle in twisted trilayer graphene (TTG) (1, 2) was soon followed by the observations of superconductivity and field-dependent quantum interference (3–5). This set of properties makes TTG the only moiré heterostructure outside of magic-angle twisted bilayer graphene (MATBG) to exhibit signatures of both a superconducting transition and macroscopic quantum phase coherence. Because TTG and MATBG share the distinctive attribute of twofold rotational symmetry $C_{2z}$, it has been proposed that this symmetry is essential to establishing superconductivity in twisted graphenes (4, 6). Superconductivity in TTG appears to be even more robust than in MATBG, with critical temperature ($T_c$) reaching up to 2.9 K in the first generation of devices. This has led to speculation that magic-angle TTG is structurally more stable than MATBG, lacking experimental devices into a mirror symmetric configuration that possesses the crucial $C_{2z}$ symmetry. Theoretical works have proposed several exotic orders for the mirror symmetric configuration, including spontaneous flavor-symmetry breaking, nematic superconductivity, and spin triplet pairing (7–9). To date, however, there is little experimental information about the atomic or electronic structure of this material; there remains no direct experimental confirmation of even the most basic hypothesis that superconducting devices possess the mirror symmetric stacking on which theoretical predictions are based.

TTG is formed by consecutively stacking three layers of graphene so that the bottom layer (B) is rotated at an angle $\theta_{BM}$ relative to the middle layer (M) and the top layer (T) is rotated at an angle $\theta_{TM}$ relative to the middle layer; both outer layers are rotated in the same direction relative to the middle layer (Fig. 1A, inset). Each rotation $\theta_i$ gives rise to a periodic density modulation, or moiré pattern, at wavelength $\lambda_d = a/\theta_i$, where $a = 0.246$ nm is the graphene lattice constant (10–12). For the special case of mirror symmetric stacking, $\theta_{BM} = \theta_{TM} = 0$ (T and B are aligned, and M is twisted relative to these by an angle $\theta$). TTG is predicted to host two sets of flat bands whose band velocity vanishes at a magic angle of $\approx 1.0$° (1, 2). As in MATBG, the quenched kinetic energy of charge carriers in these bands is expected to favor the formation of strongly correlated states of matter. Recent transport measurements have confirmed the importance of electronic correlations in TTG with the observation of superconductivity by two groups with similar phenomenology (3, 4).

Several obstacles can stand in the way of achieving perfect mirror symmetry. Despite state-of-the-art fabrication techniques, the highest-quality TTG heterostructures will inevitably have a small mismatch between $\theta_{TM}$ and $\theta_{BM}$ over macroscopic length scales, as was the case in at least one superconducting device (4). In the limit of perfectly rigid graphene layers (neglecting lattice relaxation), such a mismatch will produce a beating pattern between the top-middle (TM) and bottom-middle (BM) moirés at a “moiré of moirés” wavelength $\lambda \sim a/\delta_m$, where $\delta_m = |\theta_{TM} - \theta_{BM}|$ (Fig. 1A). In regions where the two moirés are in phase, TM AA sites sit atop BM AA sites, resulting in a locally mirror symmetric AtA (“A-twist-A”) trilayer configuration composed of AAA, ABA, and BAB stacking sites. Where the two moirés are out of phase, by contrast, the AA sites of one bilayer align with the AB sites of the other, generating a local AB stacking configuration (13, 14), comprising ABB, AAB, and BAC stacking sites; AtB is related to the AtA configuration through translation of the top layer (Fig. 1, B and C). The emergent structures of the trilayer moirés in these two regions are distinguished by their different symmetry classes, as visualized by their predicted topographic profiles in Fig. 1B. We estimate the out-of-plane corrugations for AtA and AtB domains as a superposition of sinusoidal functions of local bilayer stackings, with maxima on AA and minima on AB sites (15); we found that whereas the AtB regions host a honeycomb moiré lattice, the moiré pattern in the AtA domains is expected to be hexagonal. In this work, we used the atomic-scale imaging capabilities of ultrahigh-vacuum scanning tunneling microscopy and spectroscopy (STM/S) at temperatures from 4.8 to 7.2 K to directly characterize the electronic structure of magic-angle TTG. Our devices were fabricated by using the “cut and stack” technique, and electrical contact was made with a preplaced graphite finger to which Field’s metal µ-solder was subsequently affixed (fig. S1). STM topography of a TTG sample is shown in Fig. 1D, in which two distinct moiré wavelengths, $\lambda \sim 9$ nm and $\lambda \sim 70$ nm, are clearly visible, corresponding to the bilayer moiré and moiré of moiré length scales, respectively. The corresponding angle mismatch $\delta_m \sim a/\Lambda$ for this region is $\approx 0.2^\circ$, which is nearly identical to the mismatch of $0.3^\circ$ measured in a superconducting TTG device (4). At such small $\delta_m$, the moiré of moiré is not expected to give rise to strong direct signatures in transport, rendering microscopic probes such as STM one of the few ways of detecting it.

The presence of two moiré patterns is a generic feature over large areas of our sample (figs. S2 and S3) and represents a deviation from the three moirés ($\lambda_{TM}, \lambda_{BM}$, and $\Lambda$) that are expected on the basis of a simple rigid model (fig. S6). The STM signal in constant current mode was dominated by structural height variations across the sample surface (figs. S4 and S5), so that we could identify the global stacking configuration as AtA by the smaller moiré lattice in Fig. 1D being hexagonal rather than honeycomb at each point in space. The bright spots in topography therefore correspond to regions of local AAA stacking and were surrounded by alternating ABA and BAB domains, which we confirmed with line-cut spectroscopy (fig. S7).

The absence of AtB domains in a sample with nonzero angle mismatch $\delta_m$ implies that TTG undergoes a reconstruction on the scale...
of the moiré lattice that favors the lower-energy (1B) AtA configuration. Close examination of Fig. 1D reveals that this moiré lattice reconstruction (MLR) produces a periodic warping of the AAA site positions to enforce AtA stacking over the entire sample area. The observed warping of the moiré lattice can be understood at the atomic scale as arising from variations in the local twist angle (θA) and strain (εA) of the individual graphene layers. θA ∼ a/√la is plotted in Fig. 1, E and F, for two nearby sample regions, where la is the area of the moiré unit cell centered on position x. For small-angle mismatch δθ, the system segregates into highly uniform triangular domains (Fig. 1, E and F, blue areas) bounded by sharp point-like irregularities in the local twist angle (Fig. 1, E and F, red areas).

For Λ ≥ 30 nm, the average twist angle internal to each triangular domain (θA) saturates to a common value of ∼1.5° that is independent of the moiré of moiré wavelength (Fig. 1G).

This implies that the MLR not only enforces AtA stacking but also tends to lock the lattice to a constant local twist angle, even as θTM or θBM is varied. To shed light on this behavior, we performed structural relaxation calculations for TTG at a range of interlayer twists. The results (fig. S10) indicate that for δθ < 0.5°, θA locks to the smaller of θTM and θBM because a stronger interlayer coupling exists at a lower twist angle interface. The additional twist angle degree of freedom therefore enables TTG to locally conform to the mirror symmetric magic angle structure while “absorbing” twist angle inhomogeneity at the larger moiré of moiré length scale. The effect of the MLR on the local electronic structure is profound, as evidenced by the Fermi-level local density of states (LDOS) map (Fig. 1D, inset), which shows large modulations in the tunneling conductivity across regions of the MLR. It was therefore necessary, in considering the potentialities of TTG as a platform for correlated phases, to analyze the electronic structure on both the sub- and supra-A length scales.

In Fig. 2A, we present STM topography of a 250-nm² area, which is part of an even larger region with only a single-moiré wavelength corresponding to a twist angle of θ = 1.55°. The extreme degree of homogeneity in this area is conveyed by the local twist angle histogram (Fig. 2A, inset), showing a standard deviation of 0.03° over the entire field of view. This indicates a twist angle mismatch of δθ < 0.05°, which provided us with the opportunity to study a single domain of the MLR as well as to investigate the spectroscopic properties of a large patch of magic-angle TTG that approaches the size of a transport device.

The high energy resolution of STS permitted us to directly probe the structure of the flat bands. A series of STS measurements acquired at 7.2 K on a single AAA site is shown in Fig. 2B for a range of voltages (Vg) applied to the graphite back gate (additional twist angles are
shown in fig. S12). The measured spectrum did not change appreciably upon cooling to 4.8 K (fig. S13). At the charge neutrality point (CNP), the spectrum was dominated by a pair of overlapping resonances that arose from the partially overlapped conduction (CB) and valence (VB) flat bands. Additional soft humps at higher energy (Fig. 2B, black arrows) correspond to the edges of the next available (remote) bands. Each flat band was expected to host a saddle point in its momentum space structure, giving rise to a sharp peak, or van Hove singularity (VHS), in the density of states. We extracted the energy positions and widths of these VHSs by fitting our spectra with the sum of two Lorentzian curves and found that at CNP, the CB and VB VHSs are separated by ~18 meV and have an average width [full width at half maximum (FWHM)] of ~23 meV.

Varying \( V_g \) systematically alters the shape of the quasiparticle spectrum, changing the intensities, separations, and widths of the flat-band VHSs. In particular, we found a transfer of spectral weight between the two VHSs upon doping to the Fermi level, saturating to a minimum width of ~15 meV at \( \nu = \pm 2 \) (Fig. 2C). Last, the VHS separation is an increasing function of doping away from CNP with a distinct asymmetry between filling of electrons and holes (Fig. 2D). In general, such gate-dependent spectral shifts can be attributed either to the single-particle effect of the displacement field \( \langle D = \nabla V_g/2\hat{z} \rangle \) on the material’s band structure or to variations in the quasiparticle interaction strength as a function of band filling (\( \nu = 4k/\pi n \)), where \( \hat{z} \) is the dielectric thickness, \( n \) is the induced carrier density, and \( n \) is the carrier density at full filling of a fourfold degenerate moiré band.

We examined the role of interactions in determining the band structure of TTG by comparing the experimental spectrum with continuum model (1, 2, 17) calculations for a uniform mirror symmetric AtA stacking configuration at \( \theta = 1.55^\circ \). In Fig. 2E and fig. S14A, we compare the measured VHS separation and widths at CNP with those predicted with three separate calculations. Using inter- and intralayer tunneling parameters (18) derived from ab initio computations (19) severely underestimates both the separation and widths of the VHSs (SP1). Enhancing the monolayer graphene Fermi velocity by ~30% (SP2) enables us to reproduce the VHS separation but predicts widths that are still a factor of ~6 smaller than those found in experiment. Because a doping-dependent calculation that includes electron interactions (7) is beyond the scope of this primarily experimental work, we restricted the theoretical analysis of interactions to the CNP. Apart from spontaneous symmetry breaking, interactions can have two effects on the quasiparticle spectrum. Coulomb repulsion between electrons can change the energy landscape for a quasiparticle moving through the heterostructure, leading to a renormalization of the band structure. In addition, inelastic scattering events can lead to a finite lifetime for quasiparticle excitations, which, because of the quantum uncertainty between energy and time, causes the excitation spectrum to be broadened. Using the self-consistent Hartree-Fock procedure of (7, 18), we found that similar to the situation in MATBG (19, 20), the interaction-induced band renormalization without additional symmetry breaking accurately accounts for the separation between the peaks and roughly 70% of their widths. An additional lifetime

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**Fig. 2. Spectroscopy on a uniform 1.55° region.** (A) STM topography of a uniform area presenting a single moiré wavelength corresponding to a twist angle of 1.55°. Scale bar, 50 nm. (Top right inset) Zoomed-in topography of a single moiré unit cell showing bright AAA sites surrounded by alternating ABA and BAB domains. (Bottom left inset) Histogram of local twist angle values extracted for each moiré unit cell. Local twist angle values are as in Fig. 1, E and F. (B) AAA site STS spectra showing the evolution of the flat band structure at 1.55° as a function of applied gate voltage. Each curve represents the average of 10 measurements performed on a single AAA site from the region in (A). Gold and green arrows indicate the valence and conduction flat bands, respectively. Black arrows indicate the edges of the remote bands. Charge-neutral spectrum shows Lorentzian fits to the valence and conduction bands. Curves are offset vertically for clarity and are plotted on the same vertical scale. \( V_{\text{sat}} = 300 \text{ mV}, I_{\text{sat}} = 150 \text{ pA}, \) and \( V_{\text{mod}} = 1 \text{ mV} \). (C) FWHM of the conduction and valence band VHSs from (B) as a function of respective band filling. Each band grows flatter as it is doped to the Fermi level. (D) Separation between conduction and valence band peaks from (B) as a function of doping. (E) Comparison of VHS separation and widths at charge neutrality between experiment and three continuum model calculations. SP1 and SP2 are single-particle calculations with different inter- and intralayer hopping parameters, and HF includes electronic interactions through Hartree-Fock corrections to the continuum model, resulting in a band renormalization and lifetime broadening (18). Only the interacting calculation reproduces the experimental spectrum. (F) High-resolution AAA site spectra shifted as described in the text and with a smooth background subtracted to emphasize the evolution of the flat bands with doping. Red dashed line indicates the position of the chemical potential for which a given spectrum was acquired. Pink arrows indicate optimal doping for superconductivity. \( V_{\text{sat}} = 200 \text{ mV}, I_{\text{sat}} = 200 \text{ pA}, \) and \( V_{\text{mod}} = 0.5 \text{ mV} \).

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broadening of 4 meV is sufficient to reproduce the widths quantitatively (Fig. 2E and fig. S14A).

Tuning $V_g$ away from zero produces systematic changes in the intensities, separation, and widths of the VHSs. We examined the effect of $D$ on the single-particle band structure and found that although it does account for the shift in relative intensity of the VHSs, values of $D$ within the experimentally accessible range fail to produce notable changes in either the predicted separation or widths of the VHSs (fig. S14B). The inability of single-particle calculations to reproduce the measured quasi-particle spectrum, combined with the strong doping dependence of the latter, provides clear evidence for a pronounced band renormalization in TTG near the magic angle caused by strong quasiparticle interactions.

To isolate and better visualize this reconfiguration of the band structure, we plotted in Fig. 2F gate-dependent STS, with each curve shifted so that the flat bands remain centered on zero energy (18). The position of the chemical potential at each doping is indicated in Fig. 2F with the red dashed line. Unlike MATBG, in which superconductivity occurs in the vicinity of multiple integer filling factors of the moiré bands (21), observations of superconductivity in TTG have been strictly limited to within the vicinity of $|\nu| = 2$, with optimal doping occurring in a roughly particle-hole symmetric fashion for $2 < |\nu| < 3$ (3, 4). In MATBG, superconductivity occurs when the chemical potential is embedded in the moiré flat bands, leading to a large density of states at the Fermi level. The enhancement of the density of states relative to pristine graphene or graphite has been hypothesized to support conventional electron-phonon-mediated superconductivity (22, 23). In TTG, however, our measurements (Fig. 2F) show that large densities of states occur at multiple fillings between $\nu = -4$ and $\nu = 4$, even though superconductivity has not been observed in all of these regions in transport measurements. It is thus clear that it is not the density of states alone that controls the superconducting dome observed in transport.

One aspect of the LDOS spectrum studied in Fig. 2 is that it breaks particle-hole symmetry in a way that is not expected on the basis of noninteracting calculations. This is apparent in that the CB remains considerably broader than the VB for all measured dopings (Fig. 2C). One implication of this particle-hole asymmetry is that the chemical potential crosses the VHSs at different filling factors for electron and hole doping occurring in a roughly particle-hole symmetric fashion for $2 < |\nu| < 3$ (3, 4). However, the particle-hole asymmetry of the tunneling spectrum stands in contrast to the approximate particle-hole symmetry of the superconducting phase diagram measured in transport. This apparent discrepancy suggests that additional factors may be relevant in determining the boundaries of the superconducting phase.

Having analyzed the electronic structure at the sub-A length scale, we next turned to a detailed study of the MLR at twist angles near those for which robust superconductivity has...
been observed (3, 4). Large area topography of a sample region with angle mismatch $\delta_\theta \sim 0.25^\circ$ is shown in Fig. 3A and fig. S8A. Overlaid on the topography of Fig. 3A is a map of the local twist angle, giving a spatially averaged value of $\theta = 1.55^\circ$. The MLR segregates the system into domains of uniform moiré plaquettes arranged in a honeycomb lattice with $\theta_x \sim 1.5^\circ$, separated by quasi-one-dimensional moiré solitons. Populating the nodes of this soliton network is a hexagonal lattice of point-like faults in the local twist angle corresponding to topological moiré defects that we term “twistons.” Zoomed-in topographs of these three types of region of the MLR are shown in Fig. 3, B to D. Of the three sites, only the moiré solitons showed considerable breaking of $C_3$ rotational symmetry, which is consistent with a distinctly large value of local heterostrain $\varepsilon_x \gtrsim 0.5\%$ on these structures (24).

The large variation in $\theta_x$ and $\varepsilon_x$ on the A scale has a dramatic effect on the local electronic structure of TTG. $\theta_x$ serves as a convenient parameter to quantitatively classify

**Fig. 4. Correlated gaps and flat-band resonance.** (A to C) Gate-dependent LDOS spectroscopy on the plaquette, twiston, and moiré soliton regions. Yellow arrows in (A) and (C) indicate full filling of the moiré superlattice. Green arrows in (B) and (C) indicate correlated gaps that are confined to the twiston and soliton regions. (D and E) Continuum model calculations showing (left) the band structure and (right) density of states for twist angles of 1.45° (red) and 1.8° (blue) at two different fillings. The zero of energy corresponds to the position of the chemical potential. Flat band resonance (D) for these angles occurs when the moiré plaquette (1.45°) superlattice is filled to $n_P = 2.4$. (F) Doping-dependent LDOS spectroscopy on the twiston (blue) and plaquette (red) regions showing flat-band resonance at $|v| \sim 2.5$. Curves are offset vertically for clarity and are plotted on the same vertical scale. The vertical dashed line indicates the Fermi level. (G) Extracted values of flat-band energy splitting between twiston and plaquette sites, $\delta_c/n$ [see (E)], as a function of doping. Minima correspond to flat-band resonances and the resulting reduction in real-space electronic disorder. (H to J) Fermi-level LDOS maps at (I) charge neutrality and at the (H) electron- and (J) hole-doped flat-band resonances. Scale bars, 25 nm. $V_{\text{set}} = 300$ mV, $I_{\text{set}} = 120$ pA, and $V_{\text{mod}} = 2$ mV.
regions of the larger moiré because each region roughly corresponds to a specific value. LDOS spectra acquired on AAA sites are shown in Fig. 3E as a function of increasing $\theta_n$. Alongside these, in Fig. 3F we plot continuum model (SP2) densities of states for a series of structural parameters ($\theta$, $\epsilon$) that approximate those found in the respective experimental topography. For low relative twists corresponding to the plaquettes, the spectrum approximates that expected for TTG, with a uniform $\theta \sim 1.45^\circ$. As we increased $\theta_n$, moving onto the moiré solitons, the spectral intensity of the flat bands was progressively attenuated onto the moiré solitons, the spectral intensity decreasing uniformly way throughout all three layers. In experiments, the moiré superlattice can be identified as the carrier density at which the derivative of the chemical potential, $\partial n/\partial x$, undergoes a rapid step-like increase (Fig. 4, A and C, yellow arrows). In TTG, each moiré band is fourfold degenerate, so that full filling corresponds to a density $n_0 = 4/A$, where $A$ is the moiré unit cell area (3, 4, 29). In our case, the size of the moiré unit cell is a function of position in the MLR $[A \rightarrow V(x)]$, so that we must refer to a local filling factor $v_x = n_0(x)$. To facilitate comparisons between our local measurements and the phase diagram gleaned from bulk probe assays, we provide in fig. S8 a chart of the statistical prevalence of local filling factors as a function of induced carrier density. The area-weighted average value $v$ is an approximation of the quantity probed in transport.

In spectroscopic measurements, correlation-induced insulating states typically appear as spectral gaps centered on the Fermi level that emerge and disappear as a function of induced carrier density (30, 20, 30–32). Unlike MATBG, in which strong correlated insulating states emerge, TTG displays only weakly resistive behavior near integer fillings. It is possible that these interaction-induced resistive states (IRSs) in TTG remain relatively undeveloped because of the coexistence of an ungapped Dirac band that serves as an alternate conducting pathway (3, 4). In this scenario, we would still expect to see a suppression of the Fermi-level density of states in our spectroscopic measurements caused by the opening of an energy gap within the flat bands. In our measurements, however, we did not observe spectral gaps in uniform regions near the magic angle (Figs. 2B and 4A and fig. S13). Instead, we found that spectral gaps emerge at certain dopings near integer fillings on the twiston and soliton sites (Fig. 4, B and C, green arrows). These features of the spectrum are not expected on the basis of single-particle calculations and therefore present clear signatures of electronic correlations that are confined to particular regions of the MLR. The modulation of correlation effects by the reconstructed moiré landscape indicates the importance of the lattice reconstruction in determining the correlated phases and suggests that the microscopic structure of the MLR may have unexpected effects on bulk properties.

We next reexamined the parent state out of which superconductivity emerges, in the context of the observed MLR. The differential rates of band filling on the regions of the A-modulation (fig. S8C) mean that as we add charge to the system we are simultaneously tuning the twiston and plaquette flat bands relative both to the chemical potential and to one another. This is illustrated in Fig. 4, D and E, which shows calculated band fillings at two values of $n$ for twist angles of 1.45° and 1.8°. In Fig. 4F, we overlay the flat band spectra on twiston and plaquette sites for the full range of measured fillings. There exists a small range of $n$ for which the two sets of flat bands are maximally overlapped and in approximate resonance with one another, giving rise to an enhanced Fermi-level density of states, which favors electronic correlations. In Fig. 4G, we quantify this flat band resonance by plotting the energy difference between spatially separated flat bands as a function of doping. The resonance condition is satisfied for $2 \leq |v| \leq 3$, which is roughly aligned with the region of optimal doping for superconductivity (3, 4). We expect the range of resonant dopings to be largely independent of the particular value of the twist-angle mismatch $\delta_0$ in a given sample, given the observed relaxation phenomenon described above (Fig. 1, E to G, and fig. S10), so that the regime of optimal doping would be roughly constant across samples with $\delta_0 \leq 0.5^\circ$. Moreover, the flat band resonance occurs at dopings in between the plaquette and twiston VHSs, which is consistent with transport measurements that found superconductivity to be bounded by VHSs in doping space.

We gained further insight into the nature of the parent state by examining the effect of the flat band resonance on the real-space electronic structure through doping-dependent LDOS mapping. LDOS maps acquired at the Fermi level are shown in Fig. 4, H to J, for the three carrier densities indicated with arrows in Fig. 4G (energy dependence is provided in fig. S17). The sample displays considerable disorder at charge neutrality (Fig. 4I). The angle mismatch $\delta_0$ in this region, as in superconducting devices (4), is $\sim 0.3^\circ$, leading to magic-angle plaquettes of lateral dimension $\sim 50$ nm, which is similar in magnitude to the superconducting coherence length (3, 4). As we tuned the carrier density toward the flat band resonance, however, the LDOS maps became increasingly homogeneous (Figs. 4, H and J), indicating a reduction in the strength of the disorder potential. TTG is therefore distinct among moiré-engineered materials in that varying $V_F$ provides a means to systematically tune electronic disorder. The co-occurrence of the flat-band resonance condition, with its resulting minimization of electronic disorder and optimal doping for superconductivity, raises the possibility that the superconducting phase boundary along the doping axis is disorder.
driven. If true, this would have certain implications for the symmetry of the superconducting order parameter (33, 34). Recent transport measurements that indicate reentrant superconductivity at high magnetic field are compatible with a spin-triplet order parameter that would be sensitive to disorder of the type we observed (5).

Confirmation of this hypothesis requires direct measurements of the effect of disorder on superconductivity. Future work that systematically explores this expanded phase space by controllably tuning moiré defect density has the potential to further shed light on the pairing mechanism in TTG by determining its sensitivity to nonmagnetic impurity scattering, as has been done in a range of other unconventional systems (35–37).

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SUPPLEMENTARY MATERIALS

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Materials and Methods
Supplementary Text
Figs. S1 to S17
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Orderly disorder in magic-angle twisted trilayer graphene
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**Zooming into trilayer graphene**
Stacking and twisting graphene layers with respect to each other can lead to exotic transport effects. Recently, superconductivity was observed in graphene trilayers in which the top and bottom layers are twisted with respect to the middle layer by the same, “magic” angle. Turkel et al. used scanning tunneling microscopy to take a closer look into the stacking structure. They found that a small misalignment between the top and bottom layers caused the lattice to rearrange itself into a pattern of triangular domains. The domains had a magic-angle twisted trilayer structure and were separated by a network of line and point defects. —JS

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