Stripe phases, in which the rotational symmetry of charge density is spontaneously broken, occur in many strongly correlated systems with competing interactions\(^{1-18}\). However, identifying and studying such stripe phases remains challenging. Here we uncover stripe phases in WSe\(_2\)/WS\(_2\) moiré superlattices by combining optical anisotropy and electronic compressibility measurements. We find strong electronic anisotropy over a large doping range peaked at 1/2 filling of the moiré superlattice. The 1/2 state is incompressible and assigned to an insulating stripe crystal phase. Wide-field imaging reveals domain configurations with a preferential assignment along the high-symmetry axes of the moiré superlattice. Away from 1/2 filling, we observe additional stripe crystals at commensurate filling 1/4, 2/5 and 3/5, and compressible electronic liquid crystal states at incommensurate fillings. Our results demonstrate that two-dimensional semiconductor moiré superlattices are a highly tunable platform for which to study the stripe phases and their interplay with other symmetry breaking ground states.

Two-dimensional (2D) moiré superlattices have emerged as a powerful platform to engineer strong correlation phenomena, such as superconductivity, correlated insulators and Wigner crystal states\(^{19,20}\). With highly tunable electronic occupancy and effective interaction strength, these materials open up new opportunities to detect and study stripe phases. Nematicity (that is, electronic liquid crystal) has recently been observed in twisted bilayer graphene near integer fillings of the superlattice unit cells\(^{21,22}\). In contrast, the semiconducting transition metal dichalcogenide (TMD) moiré superlattices, whose electronic flat minibands are of distinct origins, are expected to exhibit qualitatively different correlation phenomena. The physics of the lowest-energy electronic flat miniband in TMD hetero-bilayers can be described by a single-band extended Hubbard model\(^{23,24}\) since valley degeneracy is the only remaining degeneracy. Moreover, the intrinsic lattice mismatch between the two materials gives rise to a triangular moiré superlattice even at a zero twist angle, which is also robust against small twist-angle disorders\(^{25}\). The strong light–matter interaction in 2D semiconductors further allows us to directly probe and image stripe order using high-sensitivity optical methods.

Here we report an observation of stripe phases (both incompressible and compressible) in zero-twist-angle WSe\(_2\)/WS\(_2\) hetero-bilayers at fractional fillings of the moiré superlattice cell. All devices are within ±0.5° of the zero twist angle. The near 4% lattice mismatch between WSe\(_2\) and WS\(_2\) gives a moiré periodicity of about 8 nm.

A carrier density \(n \approx 1.9 \times 10^{12} \text{ cm}^{-2}\) corresponds to filling factor \(\nu = 1\), that is, one charge per moiré superlattice cell. In the absence of a stripe phase, the in-plane linear electronic response of the system is isotropic. We probe the electronic anisotropy and compressibility of the heterostructures through resonant optical conductivity and penetration capacitance measurements, respectively, over a wide range of doping density and temperature. Details on the device fabrication and characterizations are provided in Methods.

Figure 1a,b shows the optical microscope image and side-view illustration of a representative device D1. Figure 1c shows the gate-dependent optical reflectance contrast spectrum at 5 K on the electron-doping side. The result is consistent with previous reports\(^{13,14}\). In the absence of doping (gate voltage \(<1.3\) V), the three peaks near 1.7 eV (denoted by black arrows) correspond to the intralayer moiré exciton resonances of WSe\(_2\) (ref. \(^{24}\)). The increase of the resonance amplitude at \(\nu = 1\) (gate voltage 2.25 V, denoted by green arrows) reflects the emergence of a Mott insulating state\(^{18,19}\) or a charge-transfer insulating state\(^{26}\) that is originated from a strong on-site Coulomb repulsion. The spectrum has negligible temperature dependence for the entire temperature range of study (5–40 K) (Supplementary Note 5).

We focus on the doping region between \(\nu = 0\) and 1 (enclosed by the dashed box), in which strong electronic interactions have been reported\(^{13,14}\). Figure 1d shows the filling-dependent penetration capacitance (black), which characterizes how well the heterostructure screens an applied out-of-plane electric field. Each peak corresponds to an incompressible (that is, insulating) state. We identify the reported Mott or charge-transfer insulating state at \(\nu = 1\) (magenta arrow)\(^{18,19,26}\) and the two generalized Wigner crystal states at \(\nu = 1/3\) and 2/3 originated from extended Coulomb interactions (blue arrows)\(^{14}\). We observe other insulating states, including a prominent peak at \(\nu = 1/2\) (red arrow) and weaker peaks at \(\nu = 1/4\) (purple arrow) and 2/5 (green arrow). These are also charge-ordered states from extended Coulomb interactions\(^{26}\). In Fig. 1d we compare the penetration capacitance with interlayer exciton photoluminescence (PL) intensity (blue) (see the PL spectrum in Supplementary Note 4). In the low doping regime below \(\nu = 2/3\), the insulating states show enhanced interlayer exciton PL. This property allows us to use PL, which can be measured under identical experimental conditions as the optical anisotropies, to identify the insulating states locally with a spatial resolution of roughly 800 nm (see Methods for measurement details).

We probe the electronic anisotropy in WSe\(_2\)/WS\(_2\) moiré superlattices optically. The response is crucially enhanced by choosing...
P1. The relative intensity change, $RC$, in the presence of anisotropy, from the cross-polarization configuration with that is set at angle $\phi$ from the light intensity ($R$) after an analyser (P2) (Fig. 1). A clear relative intensity change is observed to peak at $\theta = \phi \approx 0.6^\circ$ (upper panel). A clear relative intensity change is observed to peak at $\nu = 1/2$ (red dashed line) and change sign with $\phi$. Multiple control experiments, including the probe photon energy and power dependences (Supplementary Note 7), verify that the observed anisotropy is an intrinsic response of the doped electrons. The asymmetry between the two measurements is caused by the $\theta^2$ term that is not negligible when $\theta$ is comparable to $\phi$. Meanwhile the PL intensity (lower panel) also shows a prominent enhancement at $\nu = 1/2$, reflecting its insulating nature. We thus assign the 1/2 state to a stripe crystal. In contrast, the $\nu = 1/3$ and 2/3 states do not show enhanced anisotropy, although they have similar responses as the 1/2 state in the PL (blue arrows in the lower panel) and penetration capacitance. This is consistent with the previous assignment of the 1/3 and 2/3 states to isotropic electron crystals. These results are reproduced in multiple devices. Figure 2b shows the data from device D2.

We investigate the thermal melting behaviour of the stripe state at $\nu = 1/2$ in Fig. 3. We evaluate the order parameter, or equivalently $\theta$ (Fig. 3a upper panel) by antisymmetrizing the relative intensity change of Fig. 2a, $[RC(\phi) - RC(-\phi)] = 4\theta/\phi$, to remove the $\theta^2$ term. The anisotropy shows strong temperature dependence and disappears around 35 K. This is in good agreement with the temperature scale at which the penetration capacitance peak at $\nu = 1/2$ vanishes. To determine $\theta$ at each temperature more accurately, we use the optical response at the isotropic 2/3 state as a reference, that is $R(\theta = 0)$, and measure $[RC(\phi) - RC(-\phi)]$ as a function of $1/\phi$ (Fig. 3b). A linear dependence is observed for all temperatures. We extract $\theta$ from the slope. The anisotropy decreases continuously with temperature (Fig. 3c). The large uncertainty originates from...
Fig. 2 | Electronic anisotropy in WSe$_2$/WS$_2$ moiré superlattices. a, Top, doping dependence of the relative intensity change (RC) of device D1 at 5 K. The probe is at 1.685 eV. The black and grey lines are measurements at $\phi = 0.6^\circ$ and $-0.6^\circ$, respectively. Bottom, the corresponding interlayer exciton PL intensity excited by the probe light. The electronic anisotropy peaks at $\nu = 1/2$ (red dashed line), which coincides with a PL intensity peak. It is assigned to a stripe crystal state. Additional stripe crystal states are identified at commensurate fillings of 1/4 (purple), 2/5 (green) and 3/5 (orange), at which enhancement in both the electronic anisotropy and PL intensity is observed. The anisotropy signal shows marked asymmetry around $\nu = 1/2$, suggesting the importance of quantum fluctuations. b, Same as a for device D2. The probe is at 1.676 eV. Similar behaviours are observed. AU, arbitrary units.

Fig. 3 | Temperature dependence. a, Gate-dependent polarization rotation at different temperatures (upper panel) for device D1 and the corresponding penetration capacitance (lower panel). In both measurements, the peak at $\nu = 1/2$ decreases continuously with temperature and disappears around 35 K, indicating continuous thermal melting of the stripe crystal. Vertical dashed lines mark the region with optical anisotropies (upper panel) and the upper limit of the insulating region around $\nu = 1/2$ (lower panel). The optical anisotropies persist into the compressible region. The curves for different temperatures are vertically displaced for clarity. b, Dependence of the anisotropy signal as a function of analyser angle $\phi$ at $\nu = 1/2$ for several representative temperatures (symbols). The polarization rotation angle $\theta$ is obtained from the slope of the linear fittings (lines). c, Temperature dependence of $\theta$ at $\nu = 1/2$. The green dashed line marks zero rotation. Error bars are the standard deviation of $\theta$ from 3 x 3 pixels (roughly 1 x 1 μm), which is the typical scale of sample drift in the temperature dependence measurement. d, Temperature- and doping-dependent relative intensity change measured with $\phi = 0.8^\circ$. The white dashed line is a guide for the eye encircling the stripe phase region.
Fig. 4 | Stripe domains at \( \nu = 1/2 \). a, b. Spatial maps of \( \theta \) (in degrees) at two sample orientations that differ by close to 90°: \( \alpha = 16° \) (a) and 102° (b). Dashed line shows the contour of device D1. c. Sample orientation-dependent \( \theta \) at three representative points (labelled P1 to P3 in a). Symbols are experimental data; lines are fits to the theoretical relation of \( \theta = \theta_0 \sin[2(\alpha - \alpha_0)] \) (see Supplementary Note 3 for the derivation). Here \( \theta_0 \) is the anisotropy amplitude, and \( \alpha \) and \( \alpha_0 \) denote the orientation of the sample and the anisotropy fast axis, respectively. We obtain a spatial map of \( \theta \) for each sample orientation by following the design in ref. 27 to achieve both high polarization purity and spatial resolution in a wide-field microscope mode.

Next we study the stripe domains at \( \nu = 1/2 \) over the entire device D1 at 5 K in Fig. 4. To determine the anisotropy amplitude and orientation at each location, we measure the polarization rotation as a function of sample orientation and compare it to the theoretical dependence of \( \theta = \theta_0 \sin(2(\alpha - \alpha_0)) \) (see Supplementary Note 3 for the derivation). Here \( \theta_0 \) is the anisotropy amplitude, and \( \alpha \) and \( \alpha_0 \) denote the orientation of the sample and the anisotropy fast axis, respectively. We obtain a spatial map of \( \theta \) for each sample orientation by following the design in ref. 27 to achieve both high polarization purity and spatial resolution in a wide-field microscope mode.

Figure 4a shows an example at \( \alpha = 16° \). An anisotropy signal is observed in most of the regions of the device (inside the black dashed line), and its amplitude and sign vary spatially. If we rotate the sample by 90° (Fig. 4b), the sign of the signal reverses because the fast and slow axes are interchanged. Figure 4c illustrates the excellent fitting quality for three representative points on the device, labelled P1 to P3 in Fig. 4a. A map of the fitting parameter \( \theta_0 \) and \( \alpha_0 \) is shown in Fig. 4e. The colour map represents the amplitude. The length and orientation of the line segment at each position represent the local amplitude and orientation of anisotropy, respectively. Multiple domains of different stripe orientations are observed.

We analyse the frequency distribution of stripe orientation \( \alpha_0 \) over the entire device to obtain a histogram in Fig. 4d. The anisotropy (fast) axis varies over a large range of angle. The distribution shows two almost-equally large angles that differ by 90°. These angles match well the armchair and zigzag directions, respectively, of the superlattice (orange dashed lines in Fig. 4d), which are independently determined from the second-harmonic-generation measurement (Supplementary Note 2). We propose in Fig. 4f two possible charge orders for the stripe state at \( \nu = 1/2 \). The electrons form linear or zigzag stripes separated by charge-deficit regions in between. Each type of stripes has three equivalent orientations related by a 120° rotation. Most of these stripe configurations are observed with varying domain sizes and frequencies. The domain patterns change considerably after each thermal cycle, while the distribution remains peaked at the same two angles (Supplementary Note 8).

We now turn to the experimental results at filling factors away from \( \nu = 1/2 \). The anisotropy peak in doping density is substantially wider than the PL or capacitance peaks (Figs. 2 and 3a). Figure 2 shows that the enhancement in both the PL and electronic anisotropy is observed at several other filling factors \( \nu = 1/4 \) (purple dashed line), 2/5 (green) and 3/5 (orange). These are probably additional stripe crystal states (see Supplementary Note 9 for possible charge configurations). Figure 3a also shows that the anisotropy survives partially into the compressible regions of the system, where the penetration capacitance falls to the background level. (Due to the high contact resistance at low temperature, reliable capacitance data below 20 K are not available.) These finite-temperature compressible stripe states indicate the existence of electronic liquid crystal (for example, nematic and smectic) phases at incommensurate fillings. An overall phase diagram obtained by fine scans in temperature and doping at fixed \( \phi = 0.8° \) is shown in Fig. 3d. The anisotropy peak centred at \( \nu = 1/2 \) narrows with increasing temperature as the other stripe crystal and electronic liquid crystal states are melted by thermal fluctuations.
The phase diagram can be qualitatively understood from a single-band extended Hubbard model on a triangular lattice in the flat-band limit (Methods). A minimum model with up to third-neighbour interactions produces incompressible charge-ordered ground states at a sequence of commensurate fillings as found in our measurements. In particular, the ground state at \( v = 1/2 \) is an electron crystal composed of linear or zigzag charge stripes (Fig. 4f) for \( V_{r}/V_{c} < 1/2 \) and \( V_{r}/V_{c} > 1/2 \), respectively, where \( V_{r} \) is the \( r \)th nearest-neighbour Coulomb interaction. The model also predicts incompressible stripe crystal states at \( v = 2/5 \) and \( 3/5 \). The 2/5 state can be viewed as the closest packing of domain walls between the \( v = 1/3 \) and \( v = 1/2 \) domains. Extending this concept to generic fillings, the ground state at \( 1/3 < v < 2/3 \) is an incommensurate array of domain walls that are parallel to each other to avoid intersection (which costs additional energy) and is generally compressible. Such domain walls are expected to fluctuate and intersect at finite temperatures. This picture that a whole suite of stripe phases proliferate from a parental \( v = 1/2 \) stripe crystal state largely explains our observation of rotational symmetry breaking over a wide range of filling between \( v = 1/3 \) and \( v = 2/3 \). The situation is qualitatively different from that in twisted bilayer graphene\(^{2-5} \).

Above, we have only considered the electronic origin of the observed stripe phases. Our results, including the temperature and doping dependences (Fig. 3) and the overall consistency with the phase diagram of a triangular lattice Hubbard model, exclude potential trivial effects from structural and charge inhomogeneities (see Methods for details). However, they could have indirect effects on the stripe phases. For instance, the effects of uniaxial strain on stripe phases in a triangular moiré superlattice have been theoretically analysed\(^{6-9} \). Small strain increases the stripe domain size, which may account for the micrometre-scale domains observed here and in other studies\(^{10-12} \). Meanwhile, it does not strongly affect the intrinsic stripe correlations, and the stripes will preferentially align along a high-symmetry crystal axis. In contrast, large strain aligns the stripes along the strain direction and reduces the phase transition to a crossover, similar to that of a large magnetic field on ferromagnets. Unintentional strain is probably present in our samples. However, the observed alignment between stripes and crystal axes (Fig. 4d and Supplementary Fig. 10) supports the small-strain scenario. Introducing controlled strain in future studies could help to shed more light on the effects of strain on stripe phase transitions. It would be also interesting to use higher spatial resolution techniques (for example, STM) to study the stripe phases in TMD moiré superlattices.

The simple classical model discussed above neglects the quantum effects (see Methods for a discussion on its limitations). As a result, the behaviour of the \( v < 1/2 \) and \( v > 1/2 \) regions is symmetric because they are equivalent by interchanging electrons and holes\(^{6} \). The observed asymmetry of the stripe phases about \( v = 1/2 \) (Figs. 2 and 3a) indicates the importance of quantum fluctuations. Future low-temperature measurements with high-compressibility sensitivity are needed to fully understand the microscopic origin of the observed electronic liquid crystal states, particularly the role of quantum melting of stripe crystals by quantum fluctuations\(^{10-13} \). This can be investigated in 2D moiré superlattices with continuous electrostatic gating, as we have shown here. It would also be intriguing to investigate the interplay of charge stripe fluctuations with spin fluctuations, which can be conveniently probed in TMDs owing to the strong spin-orbit coupling and the spin-valley selection rules\(^{13-15} \). Our results thus demonstrate that semiconducting TMD moiré superlattices are a promising platform from which to study the cooperation and competition between the different electronic phases of matter.

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Methods

Basic properties of the WSe2/WS2 moiré superlattices. Monolayer WSe2 and WS2 have direct bandgaps at the K and K' valleys in the momentum space. The two valleys are related by time-reversal symmetry and have degenerate energy, whereas the spin-degenerate states of each valley is lifted by an external magnetic field. In the moiré superlattice, the moiré potential breaks the monolayer bands into many flat moiré minibands. The optical response is dominated by excitons, that is, tightly bound electron-hole pairs. Each optical resonance peak in Fig. 1c (labelled by black arrows) corresponds to a different exciton state in the moiré superlattice7. In the present study, we focus on the lowest conduction moiré miniband. Overtone electron doping into this band can be experimentally determined from an abrupt change in the reflectance contrast (Fig. 1c) and the penetration capacitance (Fig. 1d). Because only the valley degeneracy remains, fully occupying this flat-band requires two electrons per moiré unit cell. The filling factor ν is defined as the number of electrons per moiré unit cell.

Device fabrication. The angle-aligned WS2/WSe2 devices were made from exfoliated van der Waals materials using a layer-by-layer transfer method with a polycarbonate (PC) stamp14. All atomically thin flakes, including WSe2 and WS2, monolayers, few-layer graphite, TaSe2, and hBN, were exfoliated from bulk crystals onto Si substrates with a 300-nm oxide layer. The crystal orientations of monolayer WSe2 and WS2, flakes were determined by angle-resolved optical second-harmonic generation15–17. The flakes were aligned according to the crystal orientations during the transfer to create either a 0° or 60° twist-angle stack. The two cases can be distinguished from the second-harmonic signal of crystal orientations during the transfer to create either a 0° or 60° twist-angle stack. The two cases can be distinguished from the second-harmonic signal of crystal orientations during the transfer to create either a 0° or 60° twist-angle stack. One possibility is that stripe phases do not occur in samples with a 60° twist angle. We have ruled out trivial strain effects on the angle alignment of the WS2/WSe2 heterostructures as a function of doping density was measured in a close-cycle He cryostat (Oxford Instruments). A single-wavelength continuous-wave probe light from a Ti-sapphire laser (M Squared SOLSTIS system) was used to measure the optical birefringence of the devices (Fig. 1e). The photon energy was set to 2.683 eV (736 nm) and 1.676 eV (740 nm), respectively, for devices D1 and D2 to match their optical resonances for enhanced detection sensitivity. The light intensity used for the beating effect imaging was 1 mW to avoid saturation. To determine the beating angle ϕ, we assigned other features in both the capacitance and PL measurements to the closest rational number with a small denominator, such as 1/3, 1/2, 2/3 and so on. The PL measurement further allows us to determine the local filling factor at each position and investigate the spatial inhomogeneity in carrier density. We estimate Δν/v ∼ 10%, which is also demonstrated by the narrow peak widths of the incompressible states in PL and penetration capacitance (Figs. 2 and 3a).

Optical anisotropy measurement. A single-wavelength continuous-wave probe light from a Ti-sapphire laser (M Squared SOLSTIS system) was used to measure the optical birefringence of the devices (Fig. 1e). The photon energy was set to 1.685 eV (736 nm) and 1.676 eV (740 nm), respectively, for devices D1 and D2 to match their optical resonances for enhanced detection sensitivity. The light intensity used for the beating effect imaging was 1 mW to avoid saturation. To determine the beating angle ϕ, we assigned other features in both the capacitance and PL measurements to the closest rational number with a small denominator, such as 1/3, 1/2, 2/3 and so on. The PL measurement further allows us to determine the local filling factor at each position and investigate the spatial inhomogeneity in carrier density. We estimate Δν/v ∼ 10%, which is also demonstrated by the narrow peak widths of the incompressible states in PL and penetration capacitance (Figs. 2 and 3a).

PL measurement. Identical experimental configurations were used for the PL and the optical anisotropy measurement in Fig. 2 (above). For the PL measurement, the 800-nm short-pass filter was replaced by an 800-nm long-pass filter. Figure 2 shows the result at 12.5 K to maximize the visibility of the enhancement at filling 1/4 and 2/5. The PL spectrum in Supplementary Note 4 was measured using a different configuration. In this case, the sample was illuminated by a 352-nm excitation and the PL emission was guided into a monochromator and detected by a liquid-nitrogen-cooled CCD camera.

Sample orientation dependence. To obtain the stripe domain orientations, we measured θ as a function of sample orientation. A Sellek–Babinet compensator (Thorlabs SBC-VIS) was added right before the sample. The compensator was turned as a half-wave plate to prime the retardance at the probe wavelength, that is, to function as a half-wave plate. We rotated the compensator to vary the angle α between the incident light polarization and the sample in the xy plane. The reflected light goes through the compensator second time, and the light polarization is reverted to the original one. Rotating the compensator is equivalent to rotating the sample because the light polarization is changed only on the sample but not anywhere else.

Origins of the anisotropy. We have ruled out trivial strain effects on the observed optical anisotropy: This is possible because the doping density in moiré superlattices can be continuously tuned by electrostatic gating. The electrostatic doping accesses solely the electronic degree of freedom, and the strain-induced optical anisotropy is expected to be doping-independent. The emergence of stripe phases at specific filling factors and only at these filling factors in different devices (Fig. 2a,b) supports their electronic origin. Similarly, charge puddles formed by inhomogeneous landscape are not expected to produce such specific doping dependence. In addition, there is no strong temperature dependence and disappears around 35 K for filling factor 1/2 (Fig. 3d). It is in good agreement with the capacitance measurements for thermal melting of the charge-ordered state at 1/2 filling. In contrast, strain-induced optical anisotropy is expected to be sensitive to temperature in this temperature range.

The observation of two types of domain with anisotropy axes differed by 90° is unusual. But they are unlikely to originate from the strain effects. Particularly, Supplementary Fig. 10 shows that the anisotropy axis at a given position changes by 90° after a thermal cycle. If this were purely from strain, it would require the strain direction to change by 90°, which is unlikely especially over large sample areas and across multiple domains. Instead, a more likely scenario is that anisotropy prefers to align with one of the high-symmetry axes of the crystal (the small-strain solution) and can only change direction to change by 90°. This is unlikely especially over large sample areas and across multiple domains. Instead, a more likely scenario is that anisotropy prefers to align with one of the high-symmetry axes of the crystal (the small-strain solution) and can only change direction to change by 90°. One possibility is that the device performance is determined by the local filling factor at each position and investigated the spatial inhomogeneity in carrier density. We estimate Δν/v ∼ 10%, which is also demonstrated by the narrow peak widths of the incompressible states in PL and penetration capacitance (Figs. 2 and 3a).

Extended Hubbard model. To intuitively understand the observed stripe phases, we consider an extended Hubbard model on a triangular lattice:
The ground state is therefore solely determined by the extended inter-site Coulomb interactions, respectively. In the following analysis, we start from the first three terms in $V_{\nu}$, $V_{\nu}$, $V_{\nu}$, which correspond to the first, second and third nearest-neighbour repulsions. This model has been analysed in ref. 37 at several commensurate fillings. The ground state at $\nu = 1/3$ is an isotropic electron crystal, which is consistent with our observation. The ground state at $\nu = 1/2$ features linear or zigzag stripes depending on the ratio of $V_{\nu}/V_{\nu}$, both explicitly breaking the three-fold rotational symmetry. The model also predicts that the $\nu = 1/4, 2/5$ and $3/5$ states break rotational symmetry. The charge configurations at $\nu = 2/5$ can be obtained by tightly packing domains of the $\nu = 1/3$ state and $\nu = 1/2$ states.

Following this picture, states at a generic filling $1/3 < \nu < 1/2$ (and similarly $1/2 < \nu < 2/3$) can be obtained by adding domain walls—instead of individual particles—to the electron crystal at $\nu = 1/3$ and $1/2$. This allows us to include longer-range interactions perturbatively and estimate the compressibility of the system. At a generic filling, the ground state is an incommensurate array of domain walls $3^{3}$. The spacing between walls is determined by minimizing the total energy including the domain-wall energy and the interaction energy between walls. At doping $\nu = 1/3 + \delta$ or $1/2 - \delta > 0$, where the walls are far apart, the ground-state energy per unit length perpendicular to the wall takes the form $3^{4}$:

$$E = -\left(\frac{q}{2}\right)\delta\mu + \left(\frac{x}{2}\right)e^{-\delta l}. \tag{2}$$

Here, $l$ is the average separation of domain walls, $\delta$ is the net charge per unit length along the wall, $\mu$ is the chemical potential and $x$ describes the interaction between walls. This interaction falls off exponentially over distance, with $x$ setting the range of the screened Coulomb repulsion. Minimizing the energy yields $l = \log(1/\delta\mu)$, which varies continuously with the chemical potential. The compressibility is then given by

$$\frac{\partial n}{\partial\mu} \approx \frac{1}{\log(1/\delta\mu)} \tag{3}$$

which vanishes non-analytically at the transition from the commensurate crystal to the incommensurate domain-wall array. In this picture, rotational symmetry breaking exists throughout the filling range $1/3 < \nu < 2/3$, which provides a potential explanation of our observation.

This simple model explains the qualitative behaviour of the system observed in experiment. It is limited to the ground state and the classical limit. The quantum effects can drastically modify the stripe orders. A recent theoretical study predicted a rich quantum phase diagram for a triangular lattice Hubbard model with more realistic parameters. 3 Depending on kinetic energy and screening, multiple states can potentially emerge at $\nu = 1/2$, including the antiferromagnetic and ferromagnetic phases. Kagome lattices and tetrahedron phase, some of which break rotational symmetry. In addition, the anisotropy in incommensurate fillings can originate from fluctuations of stripes at commensurate fillings, that is, melting of stripe crystals by quantum fluctuations. 3 - 5. Metastable states such as stripe glass can also emerge. 3 - 5. Future low-temperature measurements with high-compressibility sensitivity are needed to fully understand the microscopic origin of the observed electronic liquid crystal states.

Data availability

The data that support the findings of this study are available within the paper and its Supplementary Information. Source data are provided with this paper.

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Author contributions

C.J. and Z.T. performed the optical experiments and analysis. T.L. and J.Z. performed the capacitance experiment and analysis. L.F. performed theoretical analysis. T.L., Z.T., Y.X., Y.T. and J.Z. fabricated the devices. S.L. and J.C.H. grew the bulk TMD crystals. K.W. and T.T. grew the bulk hBN crystals. C.J., J.S. and K.F.M. designed the scientific objectives and oversaw the project. C.J., J.S. and K.F.M. cowrote the manuscript. All authors discussed the results and commented on the manuscript.

Competing interests

The authors declare no competing interests.

Additional information

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