Charge-order-enhanced capacitance in semiconductor moiré superlattices

Tingxin Li¹, Jiacheng Zhu¹, Yanhao Tang¹, Kenji Watanabe², Takashi Taniguchi², Veit Elser³, Jie Shan¹,³,⁴✉ and Kin Fai Mak¹,³,⁴✉

Van der Waals moiré materials have emerged as a highly controllable platform to study electronic correlation phenomena⁴⁷. Robust correlated insulating states have recently been discovered at both integer and fractional filling factors of semiconductor moiré systems²⁰–²⁷. In this study we explored the thermodynamic properties of these states by measuring the gate capacitance of MoSe₂/WS₂ moiré superlattices. We observed a series of incompressible states for filling factors 0–8 and anomalously large capacitance in the intervening compressible regions. The anomalously large capacitance, which was nearly 60% above the device's geometrical capacitance, was most pronounced at small filling factors, below ν = 8 and anomalously large capacitance in the intervening region under a large global back gate voltage. We achieved good electrical contacts to the heterobilayer for capacitance measurements down to 10 K by using platinum electrodes³² and heavily doping the contact metal. A single-band extended Hubbard model with twofold valley pseudospin degeneracy is believed to capture the essential low-energy physics of TMD heterobilayers. In particular, the insulating state at filling factor ν = 1 (that is, one particle per moiré superlattice cell or half-filling of the first moiré miniband) is understood to originate from a strong on-site Coulomb repulsion U compared with the lattice kinetic hopping t⁹,¹². The charge-ordered insulating states at fractional fillings highlight the importance of the extended Coulomb interaction in addition to U⁹,¹². To date, the detection of these states has relied heavily on optical techniques, largely due to the challenges in achieving good TMD–metal contacts for electrical studies, particularly at low charge density and low temperature. Because of the different length scales of the Coulomb interactions that are responsible for the correlated insulating states at the integer and fractional filling factors, the degree of sample–gate coupling is expected to be different.

We report here the capacitance of field-effect devices comprising MoSe₂/WS₂ heterobilayers (Fig. 1a) as a function of filling factor spanning up to four electron moiré minibands (ν = 0–8). The angle-aligned MoSe₂/WS₂ heterobilayers form a triangular moiré superlattice with a₀ ≈ 8 nm (corresponding to a moiré density of 1.9 × 10¹² cm⁻²) because there is ~4% lattice mismatch between these materials. The MoSe₂/WS₂ heterobilayer has a type-I band alignment;¹¹ the lowest-energy electron moiré minibands are formed by states near the conduction band minimum of MoSe₂ and are weakly coupled to the states in WS₂ (Fig. 1b). We achieved good n-type contacts to the heterobilayer for capacitance measurements down to 10 K by using platinum electrodes³² and heavily doping the contact region under a large global back gate voltage Vg. The electron density in the heterobilayer can be tuned by varying the local top gate voltage Vt. We measured both the differential top gate capacitance C and the penetration field capacitance and found the two results to be fully consistent (Extended Data Fig. 1). Details of device fabrication and measurements are given in detail in the Methods.

¹School of Applied and Engineering Physics, Cornell University, Ithaca, NY, USA. ²National Institute for Materials Science, Tsukuba, Japan. ³Laboratory of Applied and Solid State Physics, Cornell University, Ithaca, NY, USA. ⁴Kavli Institute at Cornell for Nanoscale Science, Ithaca, NY, USA. ⁵E-mail: jie.shan@cornell.edu; kinfai.mak@cornell.edu
Fig. 1 | Charge-order-enhanced capacitance in MoSe2/WS2 moiré superlattices. a, Schematic of a dual-gate device structure and electrical connections for capacitance measurements. TG, top gate; BG, back gate; HEMT, high-electron-mobility transistor; \( V_{\text{meas}} \), a.c. voltage; \( V_n \) and \( I_n \), gate voltage and bias current in the HEMT, respectively; \( V_{\text{source}} \), measured voltage in the lock-in; 100 M, 100 MΩ bias resistor. b, Schematic of the moiré miniband structure of MoSe2/WS2 showing a type I band alignment. c, Schematic illustration of DOS versus energy for the first four Hubbard bands. The bandwidth increases and the DOS decreases with energy. d, Differential top gate capacitance as a function of top gate voltage at 10 K for device 1 (\( d/a_M \approx 1 \)). The back gate voltage was fixed at 5.3 V. The filling factors for discernible incompressible states are labelled on the top axis. e, Charge carrier density \( n \) as a function of top gate voltage obtained by integrating the capacitance in d in the voltage range –4.5 to –2 V. The density plateaux at incompressible states \( \nu = 1 \) and \( \nu = 2 \) are labelled.

Figure 1d shows the measured differential capacitance normalized to \( C_g \) for device 1 (\( d/a_M \approx 1 \)) as a function of top gate voltage at 10 K (see Extended Data Figs. 2 and 3 for basic characterizations of the capacitance measurements). A step increase in capacitance can be observed at around –4 V, when the Fermi level enters the conduction band of MoSe2. The capacitance plateaus out above ~0 V, when the sample is heavily electron-doped. We calibrated the capacitance using the following two limits: \( C/C_g = 0 \), when the Fermi level lies inside the superlattice bandgap (>1 eV) and the sample is insulating, and \( C/C_g \approx 1 \), when the sample is heavily doped and behaves as a good conductor. At intermediate gate voltages, we identified a series of capacitance dips (incompressible states). The most prominent ones appear equally spaced in gate voltage and have been assigned integer filling factors \( \nu = 1, 2, 3 \) and 4 (see Extended Data Fig. 4 for \( \nu \) values up to 8). These assignments are consistent with the known moiré density and the carrier density \( n \) evaluated from the gate voltage and \( C_g \approx 2.1 \times 10^{-13} \) cm\(^{-2}\). The geometrical capacitance was determined from the permittivity and thickness of the gate dielectric (hexagonal boron nitride (hBN), \( \varepsilon \approx 3\varepsilon_0 \), with \( \varepsilon_0 \) denoting the vacuum permittivity), and was independently verified by including a reference capacitor on the measurement chip (Methods).

We determined the charge density of the sample as a function of gate voltage (Fig. 1c) by integrating the differential capacitance in Fig. 1d. We calibrated the scale of \( n \) by setting it to the moiré density at \( \nu = 1 \). The incompressible states manifest a density plateau of different gate voltage width, indicating that charging into the sample is suppressed. Incompressible states at commensurate filling factors, including 1/3, 1/2, 2/3, 4/3 and 5/3, can be identified from the capacitance. The states at \( \nu < 1 \) are more robust than the states at \( \nu > 1 \). These results are in good agreement with previously reported charge-ordered insulating states in WSe2/WS2 moiré superlattices\(^{10,15}\).

Remarkably, we observed anomalously large capacitance in the compressible regions between the incompressible states. The enhancement is particularly large at small doping densities with \( C \) exceeding \( C_g \) by ~30% for device 1 (\( d/a_M \approx 1 \)). The enhancement gradually diminishes with increasing density (\( \nu \) > 2). In Fig. 2a we compare devices of different sample–gate separation (\( d/a_M \approx 0.6, 0.8, 1.0 \) and 1.5). The capacitance enhancement increases with decreasing \( d/a_M \). It is as high as 60% in device 4 (\( d/a_M \approx 0.6 \)) at 10 K. In addition, more incompressible states (marked by vertical dashed lines) can be identified between \( \nu = 0 \) and \( \nu = 1 \) in devices with larger \( d/a_M \) although they generally show less contrast due to weaker sample–gate coupling. For instance, five incompressible states (\( \nu = 1/4, 1/3, 1/2, 2/3 \) and 3/4) can be identified in the device with \( d/a_M \approx 1.5 \), but only two states (\( \nu = 1/3 \) and 2/3) are visible in the device with \( d/a_M \approx 0.6 \).

To illustrate the relationship between the observed capacitance enhancement and charge ordering, we examined the temperature dependence of capacitance for device 4 (\( d/a_M \approx 0.6 \)). Figure 3a shows only integer filling states at 100 K. On decreasing the temperature, incompressible states at fractional fillings emerge. At the same time, the capacitance in the compressible region (for example, at \( \nu = 0.22 \) and 0.52) increases below the transition temperature \( T_c \).
of the fractional states, demonstrating that enhancement is a result of the formation of charge order (Fig. 3b). Above $T_{\text{co}}$, the capacitance enhancement is small. Similar results were observed for device 1 ($d/a_d \approx 1$, Extended Data Fig. 5). The transition temperatures for the observed incompressible states are presented in Table 1.

The experimental results presented in Fig. 2a show that the correlation effects are strongly dependent on sample–gate separation, particularly for the fractional filling states because the gate electrodes effectively screen the extended Coulomb interactions. In this regime it is no longer appropriate to separate the electric potential into independent contributions of the electrostatic energy from the ideal capacitance and the sample chemical potential. Instead, we describe the entire device as a single system. The top gate voltage $V_g$ can be compared to the observed incompressible states connected by voltage jumps at the charge-ordered states. The density dependence of $V_g$ is much smaller than the nearest-neighbour Coulomb repulsion scale corresponding to the lowest measurement temperature (10 K).

Finite temperature effects alone cannot explain the discrepancy because the energy scale corresponding to the lowest measurement temperature of 10 K is much smaller than the nearest-neighbour Coulomb repulsion ($\approx 8\%$ of the total moiré potential depth). Therefore, $V_g$ is a piecewise linear function of $n$, the slopes of which are given in terms of the energies of the adjacent charge-ordered states that participate in the mixture (Methods). The density dependence of $V_g$ exhibits a series of plateaux connected by voltage jumps at the charge-ordered states. The capacitance is zero for the charge-ordered states and infinity in between.

The measured capacitance deviates substantially from this ideal limit. It is finite at the charge-ordered states: it exceeds $C_g$ by no more than 60% in the compressible regions. Finite temperature effects alone cannot explain the discrepancy because the energy scale corresponding to the lowest measurement temperature (10 K) is much smaller than the nearest-neighbour Coulomb repulsion ($\approx 12.7\%$ of the total moiré potential depth). Therefore, $V_g$ is a piecewise linear function of $n$, the slopes of which are given in terms of the energies of the adjacent charge-ordered states that participate in the mixture (Methods). The density dependence of $V_g$ exhibits a series of plateaux connected by voltage jumps at the charge-ordered states. The capacitance is zero for the charge-ordered states and infinity in between.

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shows the results obtained for different values of \(d/a_0\) at a fixed \(\varepsilon_s\) of 12.7 meV, and they are in qualitative agreement with experiment (Fig. 2a). In particular, with increasing \(d/a_0\), the capacitance enhancement in the compressible regions decreases, the incompressible dips become less prominent, but the number of robust charge-ordered states increases.

The even integer states (\(\nu = 2\) and 4) are moiré miniband insulators\(^{27,28}\), they are of single-particle origin. The odd integer states (\(\nu = 1\) and 3) are Mott or charge-transfer insulators\(^{27,28}\); they are largely driven by on-site Coulomb repulsion, which is much less sensitive to the sample–gate coupling than the extended Coulomb interaction. We thus expected the separation of the sample chemical potential \(\mu\) from the electrostatic potential of the device to be approximately valid even for \(d/a_0\approx 1\). We evaluated the sample chemical potential jump for adding or removing one electron at each integer filling state from the area of the capacitance dip, \(\Delta \mu = e \int \frac{dV}{C_{\text{cap}}}(1 - C_{\text{cap}})/2\) (Methods). By extrapolating \(\Delta \mu\) to the zero-temperature limit (Extended Data Fig. 6), we obtained the thermodynamic gap of the state (see Table 1 for device 1). We were able to reliably extract \(\Delta \mu\) from the capacitance results only for device 1, which has sufficiently small resistance at integer fillings to allow reliable capacitance measurements (Methods and Extended Data Fig. 3).

We illustrate the schematic energy diagram of the MoSe\(_2/\)WS\(_2\) moiré superlattice in Fig. 1c based on the energies inferred from the capacitance. Because the interaction effect weakens with increasing \(\nu\), we expected the Hubbard bandwidths to increase with \(\nu\). The first gap (58.5 meV, \(\nu = 1\)) is the Mott or charge-transfer gap of the first moiré miniband. The value is comparable to but smaller than the on-site Coulomb repulsion \(U\approx 200\) meV, estimated from the moiré potential amplitude from optical and scanning tunnelling microscopy measurements\(^{14-16}\) (Methods). The first moiré miniband gap (36.3 meV, \(\nu = 2\)) is smaller than the Mott gap. This indicates the influence of remote bands, that is, the mixing of different moiré minibands by \(U\) (which can also reduce the measured Mott gap). The third gap (18.1 meV, \(\nu = 3\)) is smaller than the \(\nu = 1\) gap, suggesting a smaller on-site Coulomb repulsion for the second moiré miniband than the first. This result is consistent with the larger second moiré miniband bandwidth\(^{27,28}\), dictating a larger electron

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**Table 1** | Transition temperature \(T_c\) and thermodynamic gap \(\Delta \mu\) of the incompressible states in MoSe\(_2/\)WS\(_2\) device 1 (\(d/a_0\approx 1\))

| \(\nu\) | 1/3 | 1/2 | 2/3 | 1 | 4/3 | 5/3 | 2 | 3 | 4 |
|---|---|---|---|---|---|---|---|---|---|
| \(T_c\) (K) | 40 | 40 | 40 | 160 | 30 | 30 | 75 | 80 | 25 |
| \(\Delta \mu\) (meV) | – | – | – | 58.5 | – | – | 36.5 | 18.0 | 5.0 |

The thermodynamic gaps of the integer filling states of the sample were extrapolated from the temperature dependence of \(\Delta \mu\) down to 0 K. The uncertainties in \(T_c\) and \(\Delta \mu\) were estimated to be \(\pm 5\) K and \(\pm 1.0\) meV, respectively.
Wannier orbital size and a smaller on-site Coulomb repulsion. The weakened interaction effect is also consistent with the observed $C/C_0 = 1$ at high $\nu$ (Fig. 1d).

Finally, we illustrate the extraction of the electronic entropy $S$ and specific heat capacity $T (\frac{\partial S}{\partial T})$ per moiré unit cell from the temperature-dependent capacitance in Fig. 3 for device 4 ($d/\alpha_m \approx 0.6$) and $0 < \nu < 1$ (see Methods for details). Figure 3c shows the filling factor dependence of $(S - \nu k_B \ln 2)$ at varying temperatures ($k_B$ denotes the Boltzmann constant). The entropy shows dips around the charge-ordered states at $\nu = 1/3$ and $2/3$ below $T_C$. It increases with temperature for all filling factors, but especially around $\nu = 1/3$ and $2/3$ upon melting of the charge-ordered states. This result is consistent with the electronic DOS of the moiré superlattice (the DOS of the metal gate is negligible). We also extracted the constant-density specific heat capacity $T (\frac{\partial S}{\partial T})$. Figure 3d shows broad peaks around $\nu = 1/3$ and $2/3$ close to $T_C$. The broadening of the melting transition of the charge-ordered states (instead of a sharp transition\(^1\)) presumably arises from sample inhomogeneity and disorders. The broadening of the filling factor close to $T_C$ is consistent with phase mixing of the charge-ordered states in the compressible regions; thermal melting of the phase-mixed states gives rise to heat capacity peaks away from $\nu = 1/3$ and 2/3.

In conclusion, our study has established capacitance as a powerful thermodynamic probe of correlated states in semiconductor moiré superlattices. It also illustrates the importance of sample–gate coupling and device-geometry-dependent Coulomb interactions at fractional fillings.

Online content

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References

1. Cao, Y. et al. Correlated insulator behavior at half-filling in magic-angle graphene superlattices. Nature 556, 80–84 (2018).
2. Cao, Y. et al. Unconventional superconductivity in magic-angle graphene superlattices. Nature 556, 43–50 (2018).
3. Yankowitz, M. et al. Tuning superconductivity in twisted bilayer graphene. Science 363, 1059–1064 (2019).
4. Tomarkin, S. L. et al. Electronic compressibility of magic-angle graphene superlattices. Phys. Rev. Lett. 123, 046601 (2019).
5. Chen, G. et al. Evidence of a gate-tunable Mott insulator in a trilayer graphene moiré superlattice. Nat. Phys. 15, 237–241 (2019).
6. Lu, X. et al. Superconductors, orbital magnets and correlated states in magic-angle bilayer graphene. Nature 574, 653–657 (2019).
7. Sharpe, A. L. et al. Emergent ferromagnetism near three-quarters filling in twisted bilayer graphene. Science 365, 605–608 (2019).
8. Chen, G. et al. Tunable correlated Chern insulator and ferromagnetism in a moiré superlattice. Nature 579, 56–61 (2020).
9. Serlin, M. et al. Intrinsic quantized anomalous Hall effect in a moiré heterostructure. Science 367, 900–903 (2020).
10. Regan, E. C. et al. Mott and generalized Wignier crystal states in WSe$_2$/WS$_2$ moiré superlattices. Nature 579, 359–363 (2020).
11. Tang, Y. et al. Simulation of Hubbard model physics in WSe$_2$/WS$_2$ moiré superlattices. Nature 579, 353–358 (2020).
12. Shimazaki, Y. et al. Strong correlated electrons and hybrid excitons in a moiré heterostructure. Nature 580, 472–477 (2020).
13. Wang, L. et al. Correlated electronic phases in twisted bilayer transition metal dichalcogenides. Nat. Mater. 19, 861–866 (2020).
14. Chu, Z. et al. Nanoscale imaging of correlated electronic states in WSe$_2$/WS$_2$ moiré superlattices. Phys. Rev. Lett. 125, 186803 (2020).
15. Xu, Y. et al. Correlated insulating states at fractional fillings of moiré superlattices. Nature 587, 214–218 (2020).
16. Jin, C. et al. Stripe phases in WSe$_2$/WS$_2$ moiré superlattices. Nat. Mater. https://doi.org/10.1038/s41563-021-00959-8 (2021).
17. Huang, X. et al. Correlated insulating states at fractional fillings of the WS$_2$/WSe$_2$ moiré lattices. Nat. Phys. https://doi.org/10.1038/s41567-021-01171-w (2021).
18. Ando, T., Fowler, A. B. & Stern, F. Electronic properties of two-dimensional systems. Rev. Mod. Phys. 54, 437 (1982).
19. Smith, T. P., Goldberg, B. B., Stiles, P. J. & Heiblum, M. Direct measurement of the density of states of a two-dimensional electron gas. Phys. Rev. B 32, 5609 (1985).
20. Bello, M. S., Levin, E. I., Shklovskii, B. I. & Efros, A. L. Density of localized states in the surface impurity band of a metal-insulator semiconductor structure. Sov. Phys. J. Exp. Theor. Phys. 53, 822 (1981).
21. Eisenstein, J. P., Pfeiffer, L. N. & West, K. W. Negative compressibility of interacting two-dimensional electron and quasiparticle gases. Phys. Rev. Lett. 68, 674 (1992).
22. Li, L. et al. Very large capacitance enhancement in a two-dimensional electron system. Science 332, 825–828 (2011).
23. Wu, Y. et al. Negative compressibility in graphene-terminated black phosphorus heterostructures. Phys. Rev. B 93, 035455 (2016).
24. Zondiner, U. et al. Cascade of phase transition and Dirac revivals in magic-angle graphene. Nature 582, 203–208 (2020).
25. Park, J. M., Cao, Y., Watanabe, K., Taniguchi, T. & Jarillo-Herrero, P. Flavour Hund’s coupling, Chern gaps, and charge diffusivity in moiré graphene. Nature 592, 43–48 (2021).
26. Skinner, B. & Shklovskii, B. I. Anomalously large capacitance of a plane capacitor with a two-dimensional electron gas. Phys. Rev. B 82, 155111 (2010).
27. Wu, F., Lovorn, T., Tutuc, E. & MacDonald, A. H. Hubbard model physics in transition metal dichalcogenide moiré bands. Phys. Rev. Lett. 121, 026402 (2018).
28. Zhang, Y., Yuan, N. F. Q. & Fu, L. Moiré quantum chemistry: charge transfer in transition metal dichalcogenide superlattices. Phys. Rev. B 102, 201115(R) (2020).
29. Pan, H., Wu, F. & Das Sarma, S. Quantum phase diagram of a Moire-Hubbard model. Phys. Rev. B 102, 201104(R) (2020).
30. Padhi, B., Chitra, R. & Phillips, P. W. Generalized Wigner crystallization in moiré materials. Phys. Rev. B 103, 125146 (2021).
31. Meng, Y. et al. Electrical switching between exciton dissociation to exciton funneling in MoSe$_2$/WS$_2$ heterostructure. Nat. Commun. 11, 2640 (2020).
32. Larentis, S. et al. Large effective mass and interaction-enhanced Zeeman splitting of K-valley electrons in MoSe$_2$. Phys. Rev. B 97, 201407(R) (2018).
33. Zhang, Z. et al. Flat bands in twisted bilayer transition metal dichalcogenides. Nat. Phys. 16, 1093–1096 (2020).
34. Jin, C. et al. Observation of moiré excitons in WSe$_2$/WS$_2$ heterostructure superlattices. Nature 567, 76–80 (2019).
35. Ruiz-Tijerina, D. A. & Falko, V. I. Interlayer hybridization and moiré superlattice minibands for electrons and excitons in heterobilayers of transition-metal dichalcogenides. Phys. Rev. B 99, 125424 (2019).
36. Li, H. et al. Imaging moiré flat bands in three-dimensional reconstructed WSe$_2$/WS$_2$ superlattices. Nat. Mater. https://doi.org/10.1038/s41563-021-00923-6 (2021).

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Capacitance measurement. Capacitance was measured using a close-cycle He-4 cryostat (Oxford Teslatron) connected to the sample on the same chip (several millimetres apart). The HEMT acts as a first-stage amplifier to effectively decouple the device capacitance from the parasitic capacitance from cabling. A small background capacitance remains after removing the contact resistance. The top gate voltage \( V_t \) was varied to tune the carrier density in the superlattice while the back gate voltage was fixed. Good electrical contacts to the TMD moiré superlattices were generally achieved for capacitance measurements at low doping densities (<10^{18} cm^{-3}) and low temperatures (down to ~10 K). The resultant out-of-plane electric field in the devices typically varied between ~0.3 and 0.6 V nm^{-1}. No strong electric field dependence was observed.

Calibration of capacitance results. Calibration of the capacitance results for these insulating states are non-essential because the small background does not drift in time and can be removed during the analysis. The work point of the HEMT was set by a voltage \( V_{bg} \) and a current \( I_{th} \), shown in the schematic circuit diagram in Fig. 1a. Both the top gate differential capacitance \( C(t) \) and the field penetration differential capacitance \( C_{t} \) were measured. The top gate capacitance is reported in the main text and the penetration capacitance is summarized in Extended Data Fig. 1. For the top gate capacitance measurements, we applied an a.c. voltage (10 mV in amplitude) to the TMD moiré superlattice and collected the signal from the top gate through the HEMT. The induced a.c. source-drain voltage across the HEMT was measured with a lock-in amplifier. For the field penetration capacitance measurements, the a.c. voltage was applied to the back gate and the HEMT was connected to the top gate, from which the induced charge was collected. Because of charge leakage through the bias resistor at the measurement frequency (Fig. 1a), there was a global phase shift in the a.c. signal. To determine this phase shift, we set the Fermi level inside the large heterobilayer bandgap (>1 eV). The remaining parasitic a.c. signal was pure capacitive because the sample is a good insulator and no charge carriers can be injected into the device. This allowed us to reset the global phase shift to the capacitive channel. This phase offset remained largely unchanged in the studied gate voltage range.

To exclude potential contributions from the contact resistance and/or sample resistance to the measured capacitance (especially at low doping densities and low temperatures), we studied both the in-phase (capacitive) and out-of-phase (resistive) components of the signal as a function of the frequency of the a.c. excitation voltage and sample temperature (Extended Data Figs. 2 and 3). At relatively high temperatures (>10 K), the results show that the out-of-phase component is in general negligible compared with the in-phase component from 137 to 2257 Hz. Below 10 K, the out-of-phase component gains importance with increasing measuring frequency at low doping densities (\( u_c \)). These results are not included in the main text. In addition, for devices 2–4, the resistive component at \( u_c = 1 \) and \( z = 2 \) is non-negligible up to ~20 K and down to the lowest measurement frequency due to the large contact/sample resistance. Calculations of the energy gaps from the capacitance results for these insulating states are not included. The resistive component for device 1 is negligible down to ~10 K at all filling factors (Extended Data Fig. 3). We have included in Table 1 the thermodynamic gaps evaluated for only this device. Potential differences of device 1 compared with the other devices include the presence of a small twist angle that makes the electronic bandwidth larger and the sample resistance smaller, and the presence of more in-gap states for hopping conduction.

Calibration of \( C_g \). We set the value of \( C_g \) to the measured top gate capacitance \( C(t) \) at high doping densities. We calibrated the value using a reference capacitance as described above. The calibrated value is consistent with the capacitance evaluated from the device geometry within 10% uncertainty. The uncertainty arises mainly from the uncertainty in determining the device parameters, including the thickness of the hBN gate dielectric (by atomic force microscopy), the device area (by optical microscopy) and the (out-of-plane) dielectric constant of hBN (~3). We also calibrated the value of \( C_g \) using the charge carrier density versus voltage curve in Fig. 1e. We set the density at the \( z = 1 \) plateau to the moiré density (1.9 × 10^{19} cm^{-2} with an estimated uncertainty of 2%) and determined \( C_g \) from the measured gate voltage. All the values are consistent with each other.

Estimation of \( U \). We estimated the magnitude of the on-site Coulomb repulsion for the first electron moiré miniband in MoSe_{2}/WS_{2} moiré superlattices using \( U = \mu e \). Here, \( \mu \approx 4 \times 10^{3} \) eV cm^{-2} with an effective permittivity of the background (hBN) and \( \xi \) is the size of the lowest-energy local Wannier orbital. We estimated \( \xi = \sqrt{\pi} \), \( \xi = 2 \) by approximating the moiré potential as a harmonic potential around its minimum. Here, \( a_{bg} (\approx 8 \text{ nm}) \) is the moiré period, \( m (\approx 0.5 \text{ nm}) \) is the conduction band mass of monolayer MoSe_{2}, and \( V_{SL} \) is the amplitude of the moiré potential (the potential depth is up to 6\( V_{SL} \) in this definition). The constants \( e, \kappa, m \) and \( h \) denote, respectively, the elementary charge, the vacuum permittivity, the free electron mass and the reduced Planck’s constant.

Lumped circuit model for the capacitance measurement. We considered a geometry in which the sample is equally spaced from the top to the back gates by distance \( d \). A lumped circuit model for the capacitance measurement is appropriate (Extended Data Fig. 8) when the interaction between electrons in the sample does not depend on the device geometry and the electrochemical potential on the gate can be separated into independent contributions of the electrostatic potential from the ideal capacitance and the sample chemical potential. We derived from the circuit the top gate capacitance

\[
\frac{1}{C_{t}} \approx \frac{2}{C_{\xi}} + \frac{1}{C_{g}} \tag{1}
\]

and for the penetration capacitance \( C_{\xi} \)

\[
\frac{1}{C_{\xi}} \approx \frac{2}{C_{\xi}} + \frac{C_{Q}}{C_{\xi}} \tag{2}
\]

Here, \( C_{Q} \) denotes the sample quantum capacitance and \( C_{\xi} \) is the geometrical capacitance between the sample and the top (or back) gate electrode. Because the a.c. excitation voltage is applied to the sample and the sample is nearly equally spaced from the top to the back gates, the charge modulation on the sample is about twice the voltage that is collected from the top gate. This accounts for the factor of two in equations (1) and (2). All capacitances are expressed in capacitance per unit area. An incompressible state appears as a dip in the top gate capacitance, and as a peak in the penetration capacitance, as a function of gate voltage.

As discussed in the main text, this approximation is appropriate for integer filling states. The chemical potential jump \( \Delta \mu \) at an integer filling state can be obtained from the gate dependence of the top gate capacitance as

\[
\Delta \mu = \int \frac{d\mu}{dV} = \frac{1}{2} \int \frac{dV}{C(t)} \left( 1 - \frac{C_{\xi}}{C_{g}} \right) \tag{3}
\]

and the penetration capacitance as

\[
\Delta \mu = \int \frac{d\mu}{dV} = \int \frac{dV}{C_{\xi}} \frac{C_{Q}}{C_{\xi}} \tag{4}
\]

The integration range includes the corresponding capacitance dip/peak in gate voltage. The value of \( \Delta \mu \) is half of the dip area from the top gate capacitance measurement, and the full peak area from the penetration capacitance measurement.

General top gate capacitance model. When the interaction between the charges in the sample is strongly dependent on the device geometry, it is no longer appropriate to separate \( eV \) into independent contributions and apply the lumped circuit model\(^1\). Therefore, we derived the capacitance as a function of the energy density \( u \) of the entire device. In general, the top gate has electron density \( n_t \), the back gate has electron density \( n_{bg} \), and the sample has electron density \( n \). Defining \( \Delta n = (n_t - n_{bg})/2 \), we can express the gate densities as \( n_t = n_{bg} = n - \Delta n \) and \( -n_{bg} = n - \Delta n \) and write the energy density of the entire device as

\[
u(n_t, n_{bg}) = u(n_t) + \frac{(e\Delta n)^2}{C_g} \quad \tag{5}
\]

where \( u(n) \) is the energy density of the sample when it holds electron density \( n \) and is surrounded by gates with equal, compensating densities \(-n/2\). The separation of the energy into the two independent contributions in equation (5) follows from the superposition of the corresponding electric fields. The field produced by \( e \) is symmetric about the sample plane whereas the electric field produced by \( e\Delta n \) on the two gates is antisymmetric with respect to that plane, with the result that there is no cross-term in the electric energy density.

The electron densities \( n_t - n \) and \( n - n_{bg} \) of the top and back gates, respectively, are kept in equilibrium by voltages \(-V_t\) and \(-V_{bg}\) on these gates, relative to the sample, provided

\[
V_t = \frac{1}{e} \frac{\partial u}{\partial n_t}, \quad V_{bg} = \frac{1}{e} \frac{\partial u}{\partial n_{bg}} \quad \tag{6}
\]
Using equation (3), these equations are equivalent to
\[
\frac{V_1 + V_0}{2} = \frac{n_i^e (n_i + n_j)}{e}
\]
\[
\frac{V_1 - V_0}{2} = \frac{\epsilon (n_i - n_j)}{2C_0},
\]
where \(\epsilon = \frac{d}{\text{area}}\). Solving equation (8) for \(n_j\) and substituting into equation (7) we obtain
\[
\frac{V_1 + V_0}{2} = \epsilon [2m_0 - C_0 \left( V_1 - V_0 \right) / \epsilon] / e.
\]
In general, the density response with respect to arbitrary variations in \(V_1\) and \(V_0\) can be obtained from this equation. We measured the top gate differential capacitance \(C = e^2 / d\) under the constraint \(dV_i = dV_0\) and \(dn_i = dn_0 = dn/2\) because the same a.c. voltage \(dV_i\) is applied between the sample and the two gates. Applying this to equation (9) we obtain
\[
C^{-1} = 2\epsilon^2 / e^2.
\]
Following a similar procedure, we can also derive the penetration capacitance \(C^p = \epsilon d / 4n_0^e\) (with \(dV_0 = 0\) as
\[
\tilde{C} = \frac{C^p}{\epsilon} = 2\epsilon d / 4n_0^e.
\]
Microscopic model for \((u_n)\). Geometrically, we modelled the device as a thin sample of area \(A\) between the top and back gate capacitances equally spaced by distance \(d\). The region between the conductors is filled with a uniform dielectric medium of permittivity \(\epsilon\). The moiré superlattice constant is \(a_{\text{BL}}\) and the unit cell area is \(A_{\text{BL}} = \sqrt{\sqrt{d}} A / 2\).

We analysed the zero-temperature limit of the classical Hamiltonian
\[
H = \frac{1}{2} \sum_{i,j} V(r_{ij}) \epsilon n_i n_j + \sum_i D_i n_i,
\]
where \(n_i \in \{0, 1\}\) is the occupation of moiré site \(i, D_i\) is the on-site disorder energy at that site, \(r_{ij}\) is the distance between sites \(i\) and \(j\), and
\[
V(r, d) = \frac{\epsilon^2}{4\pi A d^3} \sum_{k=\text{odd}}^{\infty} \frac{(-1)^k}{\sqrt{k^2 + (2kd)^2}}
\]
is the electrostatic energy in the symmetrical two-conductor geometry (\(k\) is an integer). The ratio \(\epsilon = H / A\) is the sample energy density defined above. All of our energies, when not explicitly specified, are reported in the energy unit \(\epsilon_a = \epsilon^2 / 4\pi A d^3 \approx 50\text{ meV}\).

The values of \(D_i\) are distributed in the range \([-\delta, \delta]\) with the distribution function \(\delta\) is a dimensionless disorder potential strength
\[
\rho(\epsilon) = D(\delta - e^2)^2 / 2,
\]
where the normalization constant \(D\) was chosen so that the integral of the distribution function \(\rho(\epsilon)\) over the dimensionless energy \(\epsilon\) is 1. For this disorder distribution, the parameter \(\delta = 0.7\) corresponds to an RMS disorder strength of 12.7 meV.

In the weak-disorder limit we may use the following simple model for the ground states of \(H\). Assuming that \(D\) varies smoothly on the moiré superlattice, the ground states are coexisting charge-ordered states with domains selected such that the densest state occupies the sites with lowest \(D_i\), followed by the next-densest charge-ordered state in the next-lowest sites, and so on. Let \(e_i > e_j > \cdots > e_{k+1}\) be the filling factors of up to \(m + 1\) charge-ordered states that participate when the average filling is \(\nu\), and the associated site-energy boundaries are
\[-\delta = e_0 < e_1 < \cdots < e_m < e_{m+1} = \delta.\]
The constraint that all electrons are in a charge-ordered state with the correct (decreasing) order of fillings is expressed as
\[
\nu = (e_1, \cdots, e_m) = \sum_{k=1}^{m+1} e_k \int_{e_{k-1}}^{e_k} \rho(\epsilon) \, d\epsilon.
\]
where \(\mu_k = \frac{1}{2} \sum_{k=1}^{m+1} V(r_{ij}) \epsilon n_i n_j\) is the electrostatic energy of an electron at site 0 in the \(k\)th charge-ordered state with occupations \(n(k)\), for which the expression for \(\mu_k\) needs to be averaged over site 0 when the electrons in the charge-ordered state are not equivalent by symmetry (for example, \(\nu = 3/3\)). The chemical potentials, labelled according to their filling factor, \(\mu_k\), are given in Extended Data Fig. 9 (in units of eV) for the fillings we identified in our earlier experiment \(1^\text{st}\) with \(d/a_{\text{BL}} = 5\), but now for the smaller values of \(d/a_{\text{BL}}\) relevant to the current experiment. Thermodynamic stability in the zero-disorder limit requires that the chemical potential \(\mu\) of a pure phase with filling factor \(\nu\) is lower than the energy of a mixture formed by the adjacent phases with fillings \(\nu_1\) and \(\nu_2\), and the corresponding \(\mu_1\) and \(\mu_2\). This translates to the inequality
\[
(\nu_1 - \nu) \epsilon < (\nu_2 - \nu) \epsilon + (\nu - \nu_1) \epsilon \mu_1 + (\nu - \nu_2) \epsilon \mu_2,
\]
which is satisfied for all the filling factors and the \(d/a_{\text{BL}}\) in Extended Data Fig. 9.

The ground-state energy per unit area \(u_a(n)\), for average filling \(\nu\), is defined as the solution of a constrained minimization problem:
\[
A_{\text{BL}} u_a(n) / A_{\text{BL}} = \min_{\nu_1, \cdots, \nu_m, \cdots, \nu_{m+1}, \cdots} u_a (\nu_1, \cdots, \nu_m) \cdot \ln \nu_m / \ln \nu_{m+1}.
\]
We evaluated \(u_a\) for our model as a function of \(d/a_{\text{BL}}\) for fillings in the range \(0 < \nu < 1\) and for several strengths of disorder \(\delta\). In the clean limit \(\delta = 0\), \(u_a\) is a piecewise linear function of density, given by weighting the pure-phase energies of the adjacent charge-ordered states. As discussed in the main text, the capacitance is zero for pure phases and infinity for mixtures.

With finite disorder \(\delta > 0\), the constrained minimization in equation (20) was solved numerically on a fine grid of \(\nu\) values using the FindMinimum function of Mathematica (https://reference.wolfram.com/language/ref/FindMinimum.html), and the second derivative \(u''_a\) was computed by finite differences. Selected results are shown in Fig. 2b,c. In addition to the incompressible \(\nu = 1\) state, we also observed incompressible behaviour at \(\nu = 0, 1/3, 2/3, 1/2\), and to a smaller degree, at \(\nu = 1/2\). Incompressibility at \(\nu = 0\) is strictly a consequence of disorder. The pockets of \(\nu = 1/7\) domains in the \(\nu = 0\) limit of our model are just individual electrons optimally occupying the sites with the lowest disorder potential. If disorder played no role in the experiment, then the resulting incommensurate Wigner crystal for \(\nu = 0\) would exhibit diverging compressibility, and not the vanishing compressibility we observed. In addition, because the \(\nu = 1/2\) stripe crystal state is driven by competing long- and short-range Coulomb repulsions, it is more sensitive to disorder and local strain than the \(\nu = 1/3\) and \(\nu = 2/3\) states, especially at smaller values of \(d/a_{\text{BL}}\). We observed this same behaviour in the experiment in which we compared the two measurements at 10 K in Fig. 2a. For the same reason, a stronger sample-to-sample variation was also observed for the \(\nu = 1/2\) state.

Device entropy. At finite temperature, the energy density \(\epsilon\) in the above treatment is replaced by the free-energy density \(f\). Equation (6) now becomes
\[
\nu_1 = 2 / (\epsilon / \text{area})^2 \nu_1.
\]
where \(\nu_1 < \nu_2 < \cdots < \nu_m < \cdots < \nu_{m+1} < \cdots\)
\[
S = \int d\nu \left( \frac{\partial S}{\partial \nu} \right)_\nu = -\frac{1}{dA_{\text{BL}}} \int dV_i \left( \frac{\partial C}{\partial \nu} \right)_\nu.
\]
The temperature dependence of the top gate capacitance \(C\) from experiment is shown in Fig. 3b for several representative filling factors of device 4 (\(d/a_{\text{BL}} \approx 0.6\)). All the curves are smooth. We could therefore first fit the experimental data with a polynomial function (dashed lines in Fig. 3b) and then evaluate the temperature derivatives using the fitted functions. The entropy extracted from experiment consists of contributions other than the electronic entropy. For instance, the MoSe₂/WS₂ bandgap is temperature-dependent. The \(\delta\) threshold corresponding to the onset of doping in the moiré superlattice exhibits a systemic shift with temperature. The main effect of such a shift is to add a constant to \(\frac{\partial S}{\partial \nu} \nu\), or a linear filling dependence to the entropy. In our analysis for temperatures ranging from 20 to 60 K, we subtracted a constant from \(\frac{\partial S}{\partial \nu} \nu\), so that \(S = 0 \text{ at } \nu = 0\) and \(S/k_B \approx 2n\text{ at } \nu = 1\) (to account for the spin-valley entropy). This is well justified because at \(\nu = 0\) the thermally excited carriers are negligible due to the large energy gap in MoSe₂/WS₂ (\(\sim 1\) eV). In the case of half-band filling (\(\nu = 1\)), the Coulomb gap is much bigger than both.
the thermal energy and the moiré bandwidth (<10 meV) so that the doublon density is negligible; the exchange interaction energy is much smaller than $k_B T$ (ref. 11) so that the spins/valleys at different moiré sites are independent. However, between $\nu = 0$ and $\nu = 1$ empty sites become available for electron hopping without creating doublons. The available microstates for hopping are expected to be largest at around half-filling of the lattice ($\nu = 1/2$). This electronic motional degree of freedom can be treated as independent of the spin/valley degree of freedom when $k_B T$ is much larger than the exchange interaction energy. We can therefore subtract $\nu k_B \ln 2$ from $S$ to isolate the electronic entropy from the motional degree of freedom. In Fig. 3c, we show $(S - \nu k_B \ln 2)$ to highlight this effect.

**Data availability**

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

**References**

37. Wang, L. et al. One-dimensional electrical contact to a two-dimensional material. Science **342**, 614–617 (2013).
38. Ashoori, R. C. et al. Single-electron capacitance spectroscopy of discrete quantum levels. Phys. Rev. Lett. **68**, 3088 (1992).
39. Zibrov, A. A. et al. Tunable interacting composite fermion phases in a half-filled bilayer graphene Landau level. Nature **549**, 360–364 (2017).
40. Shi, Q. et al. Odd- and even-denominator fractional quantum Hall states in monolayer WSe$_2$. Nat. Nanotechnol. **15**, 569–573 (2020).
41. Saito, Y. et al. Isospin Pomeranchuk effect in twisted bilayer graphene. Nature **592**, 220–224 (2021).
42. Rozen, A. et al. Entropic evidence for a Pomeranchuk effect in magic-angle graphene. Nature **592**, 214–219 (2021).

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**Author contributions**

T.L. performed the measurements. T.L. and J.Z. fabricated the devices and analysed the data. Y.T. provided assistance in the device fabrication. V.E. performed theoretical calculations. K.W. and T.T. grew the bulk hBN crystals. T.L., J.S. and K.F.M. designed the scientific objectives and oversaw the project. T.L., V.E., J.S. and K.F.M co-wrote 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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**Correspondence and requests for materials** should be addressed to J.S. or K.F.M.

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Extended Data Fig. 1 | Penetration capacitance. Penetration capacitance (a) and top gate capacitance $C/C_g$ (b) as a function of top gate voltage at 20 K ($d/a_m \approx 0.6$). We also show the top gate capacitance (obtained from a) in b for comparison. Good agreement between the penetration and top gate capacitance measurements is observed.
Extended Data Fig. 2 | Frequency dependent capacitance ($d/a_{M} \approx 0.6$). Top gate capacitance as a function of top gate voltage at 10 K (a) and at 20 K (b) under different excitation frequencies (10 mV modulation amplitude). A decrease in capacitance with increasing frequency is observed at 10 K, signifying the growing importance of the resistive component in the AC measurement. Meanwhile, no obvious frequency dependence is seen at 20 K, a temperature at which the sample/contact resistance becomes substantially smaller. The capacitance measurement is therefore in the low-frequency limit and reliable for most of the filling factors except $\nu = 1$ and $\nu = 2$. 
Extended Data Fig. 3 | Capacitive and resistive channels. a, Top gate (or filling) dependence of the capacitive channel output (top) and the resistive channel output (bottom) at different excitation frequencies (device 4 at 10 K). A constant background was subtracted from the capacitive channel for better comparison. The resistive signal gradually goes away with decreasing frequency for most of the filling factors except \( \nu = 1 \) and \( \nu = 2 \), where a small increase instead is observed. We conclude that the AC capacitance measurement at 10 K is reliable for most of the filling factors at the lowest frequency (137 Hz) for this device. However, for \( \nu = 2 \) and \( \nu = 1 \), the AC measurement is in the high-frequency limit so that the measured capacitance dips at \( \nu = 2 \) and \( \nu = 1 \) only reflect the large sample/contact resistance. b, Same as a measured at different temperatures (device 1 at 3137 Hz). As the resistive signal grows with decreasing temperature, the capacitive signal decreases accordingly at 10 K and above. The resistive signal is negligible compared to the capacitive signal. In addition, small resistive peaks at \( \nu = 2 \) and \( \nu = 1 \) grow with decreasing temperature. All of these observations show that the measurements are in the low-frequency limit for all filling factors in this device at 10 K and above.
Extended Data Fig. 4 | Top gate capacitance of device 2. The top gate capacitance shows incompressible states up to $\nu = 8$. 
Extended Data Fig. 5 | Temperature dependence of device 1. 

**a**, Top gate capacitance as a function of filling factor at different temperatures. The charge-ordered states melt and the large capacitance enhancement disappears at high temperatures.  

**b**, Top gate capacitance at selected filling factors as a function of temperature. The capacitance in the incompressible region ($\nu = 1/3$ and $1/2$) decreases rapidly below their transition temperatures, whereas the capacitance in the compressible regions ($\nu = 0.42$) increases.  

**c**, Filling factor versus top gate voltage at different temperatures obtained from integrating the data in **a**. The inset shows the temperature dependence of the top gate voltage at two constant filling factors.  

**d**, Filling factor dependence of $(S - k_B \ln 2)$, where $S$ is the electronic entropy per moiré unit cell of the device extracted from two temperature windows, 15-40 (black) and 40-80 K (red).
Extended Data Fig. 6 | Temperature dependence of $\Delta \mu$ (device 1). Thermodynamic gap size $\Delta \mu$ as a function of temperature for $\nu=1$ to 4 extracted from the capacitance measurements. Linear extrapolation (dashed lines) was used to extract the zero-temperature gap size tabulated in Table 1.
Extended Data Fig. 7 | Optical micrograph. Optical image of a typical device.
Extended Data Fig. 8 | Lumped circuit model. Lumped circuit model for the top gate capacitance measurement (a) and penetration capacitance measurement (b).
Extended Data Fig. 9 | Table for the calculated chemical potentials. Calculated chemical potentials, $\mu_\nu$ (in units of $\epsilon_a$) at different filling factors $\nu$ and $d/a_M$.

| $\nu$  | $d/a_M = 0.75$ | $d/a_M = 1.00$ | $d/a_M = 1.25$ | $d/a_M = 1.50$ |
|-------|----------------|----------------|----------------|----------------|
| 1/7   | 0.017          | 0.059          | 0.126          | 0.211          |
| 1/4   | 0.076          | 0.198          | 0.357          | 0.538          |
| 1/3   | 0.146          | 0.339          | 0.573          | 0.828          |
| 2/5   | 0.269          | 0.525          | 0.821          | 1.137          |
| 1/2   | 0.409          | 0.754          | 1.140          | 1.547          |
| 3/5   | 0.538          | 0.973          | 1.449          | 1.947          |
| 2/3   | 0.611          | 1.103          | 1.639          | 2.197          |
| 3/4   | 0.743          | 1.311          | 1.922          | 2.556          |
| 6/7   | 0.900          | 1.566          | 2.275          | 3.006          |
| 1     | 1.076          | 1.867          | 2.705          | 3.566          |