Imaging field-tuned quantum Hall broken-symmetry orders and the quantum Hall conducting channel in a charge-neutral graphene/WSe$_2$ heterostructure

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The zeroth Landau level (0LL) in graphene has emerged as a flat band platform in which distinct many-body phases can be explored with unprecedented control by simply tuning the strength and/or direction of the magnetic field. A rich set of quantum Hall ferromagnetic phases with different lattice-scale symmetry-breaking orders are predicted to be realized in high magnetic fields when the 0LL in graphene is half-filled. Here we report on a field-tuned continuous phase transition of different valley orderings in a quantum Hall ferromagnetic phase of charge-neutral graphene on insulating tungsten diselenide (WSe$_2$). The phase transition is clearly revealed by an anomalous field-dependent energy gap in the half-filled 0LL. Using atomic resolution imaging of electronic wavefunctions during the phase transition, we unexpectedly observe the microscopic signatures of field-tuned continuous-varied valley polarization and valley inversion, which are beyond current theoretical predictions. Moreover, the quantum Hall conducting channel of the graphene is directly imaged when the substrate (WSe$_2$) introduces band bending of the 0LL.

graphene, WSe$_2$, quantum Hall ferromagnetic, valley polarization, valley inversion

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1 Introduction

In a flat band, the kinetic energy is quenched, and electron-electron interactions can drive electrons to form exotic correlated phases to minimize the Coulomb interaction. For example, correlated insulator, superconductivity, and the quantum anomalous Hall effect were observed very recently in partially filled flat bands of magic-angle twisted bilayer graphene [1-5]. In a graphene system, introducing a perpendicular magnetic field is the simplest way to create flat bands, i.e., Landau levels (LLs), and realize exotic many-body phases [6-28]. At the half-filled zeroth LL (0LL), the Coulomb interaction gives rise to a particular SU(4) quantum Hall ferromagnet (QHFM) because of the fourfold spin-valley invariance symmetry in a graphene monolayer. Then, the spin and valley anisotropic energy terms explicitly break the SU(4) symmetry and lead to a diversity of quantum Hall ferromagnetic phases with distinct broken spin and valley symmetry [6-13]. Although previous studies have achieved great success in exploring the QHFM [14-17,22-27], the nature of the ground states of charge-neutral graphene, which is determined by a delicate balance between the spin and valley anisotropic energy terms, has remained under intense debate.

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In the 0LL of graphene, the wavefunctions of each valley are locked to one of the real space sublattices, enabling us to precisely pinpoint the nature of the quantum Hall ferromagnetic phases by directly imaging the atomic-scale wavefunctions in real space, as first demonstrated in ref. [24] and further confirmed very recently [26,27]. In this work, using atomic resolution imaging of electronic wavefunctions of charge-neutral graphene on insulating tungsten diselenide (WSe$_2$), we report the microscopic signatures of a field-tuned continuous phase transition of different valley orderings in the QHFM. During the phase transition, we observed field-dependent valley polarization and inversion, which are closely related to the observed anomalous field-dependent energy gap between 11.5 and 14 T in the half-filled 0LL. Our work demonstrates the power of exploring the nature and phase transition of many-body orders by real-space atomic-scale imaging of the wavefunctions. The unexpected results obtained in this work, as well as the complex orders reported in refs. [14-16,24-27], indicate that the true phase diagram of the QHFM of charge-neutral graphene is probably sample- and/or substrate-dependent, maybe partially arising from their different field-dependent spin and valley anisotropic energy terms, and needs further investigation.

2 Experimental results

In our experiment, the graphene/WSe$_2$ heterostructure was obtained using a wet transfer fabrication of single-crystal monolayer graphene on mechanical-exfoliated thick WSe$_2$ sheets (using the transfer platform from Shanghai Onway Technology Co., Ltd., see Supplementary Materials Figure S1(a) for details) [29,30]. Figure 1(a) shows a typical scanning tunneling microscope (STM) image of the graphene monolayer on a multilayer WSe$_2$ substrate, and the inset shows its corresponding fast Fourier transforms (FFT) image. We can observe the moiré superlattice structure due to the lattice mismatch between WSe$_2$ and graphene (lattice constant: 0.335 nm for WSe$_2$ and 0.246 nm for graphene). Our experiment indicates that the Dirac point of graphene in the graphene/WSe$_2$ heterostructure depends on the thickness of WSe$_2$, and the graphene can be changed from p-type doping to n-type doping using the thickness of WSe$_2$ (Figure S1(b) and (c)). Therefore, we can study the electronic properties of charge-neutral graphene by carefully choosing the thickness of the supporting WSe$_2$. Figure 1(b) shows scanning tunneling spectroscopy (STS) measurements of two representative graphene/WSe$_2$ heterostructures labeled Device 1 and Device 2 in a perpendicular magnetic field $B=10$ T. The tunneling spectra of both devices exhibit well-defined Landau quantization of massless Dirac fermions (Figures S2 and S3), as expected from a graphene monolayer. The doping of the two devices is very different. The 0LL is approximately 35 meV above the Fermi level for Device 2 but exactly at the Fermi level, i.e., half-filled, for Device 1. A pronounced feature is the splitting of the half-filled 0LL with a gap of $\Delta E \approx 24$ meV in Device 1, indicating a Coulomb gap when a flat band or (quasi)bound state is partially filled [24,26,27,31-36]. The most important feature of a Coulomb gap of the 0LL is its dependence on the filling of the 0LL: We can observe the Coulomb gap only when the 0LL is partially filled. This result can help us exclude the possibilities of an insulating gap generated by the sample’s disorder or possible graphene-substrate coupling. Moreover, the observed gap agrees quite well with the on-site Coulomb energy of $E_C = e^2 / 4\pi\varepsilon \varepsilon_0 l_B \sim 25$ meV, where $e$ is the electron charge, $\varepsilon_0$ is the vacuum dielectric constant, $\varepsilon_r = (1 + \varepsilon_{WSe_2}) / 2 \approx 4.4$ is the relative dielectric constant surrounding the graphene ($\varepsilon_{WSe_2} \approx 7.8$ is the relative dielectric constant of WSe$_2$ [37,38]), and $l_B = \sqrt{\hbar / eB}$ is the magnetic length, with $\hbar$ the reduced Planck constant. This agreement indicates that the Coulomb interaction lifts the degeneracy of the half-filled 0LL, and the gain in the exchange energy may favor the formation of spin/valley polarized/coherent orders in charge-neutral graphene.

To further study the Coulomb gap of the 0LL, we measured high-resolution tunneling spectra in different magnetic fields in Device 1 (Supplementary Materials for details). Figure 1(c) shows the 0LL as a function of the magnetic field $B$. In the small magnetic field region ($B<8.2$ T), the 0LL is filled (with a filling factor of 2), and there is no Coulomb gap, further confirming that the splitting of the half-filled 0LL arises from the Coulomb interaction. The position of the 0LL depends strongly on the magnetic field for $B<8.2$ T because of competition between magnetic confinement and spatial confinement generated by the tip-induced electrostatic potential [39,40]. When $B>8.2$ T, the 0LL is half-filled, and the position of the 0LL is independent of the magnetic field, indicating that the magnetic length $l_B$ is much smaller than the characteristic length of the tip-induced potential, and the effect of the tip gating is negligible [39]. Then, the observed gap at the half-filled 0LL should be described by the on-site Coulomb energy. Figure 1(d) summarizes the field-dependent splitting of the 0LL obtained in Device 1. The filling factor is zero for $B>8.2$ T, and the splitting of the 0LL for $8.2$ T $< B < 11.5$ T agrees quite well with the on-site Coulomb energy $E_C = e^2 / 4\pi\varepsilon \varepsilon_0 l_B$ of the graphene monolayer on the WSe$_2$ substrate. However, the observed gap for the half-filled 0LL obviously deviates from the on-site Coulomb energy for $B>11.5$ T. For simplicity, the observed charge gap in the 0LL of graphene should be the sum of the exchange (Coulomb) term and the lattice-scale symmetry-breaking terms, both of which can be strongly modified by the strength and/or direction of the magnetic field [5-8]. Therefore, the
observed anomalous field-dependent charge gap in the half-filled 0LL for \( B > 11.5 \, \text{T} \) suggests field-induced phase transitions between distinct lattice-scale orders in graphene.

To further explore the origin of the anomalous field-dependent charge gap in the half-filled 0LL, we use STS maps of the occupied and empty states of the 0LL at the atomic scale. Although we observe well-defined hexagonal honeycomb lattices at the Dirac point of graphene in a zero magnetic field (Figure S4), the STS maps of the occupied and empty states of the 0LL at \( B = 10 \, \text{T} \) exhibit a triangular lattice density distribution (Figure 2(a), left panels). This result indicates that the ground state of the charge-neutral graphene on WSe\(_2\) at 10 T is the valley-polarized charge density wave (CDW) phase according to the sublattice-valley locking in the 0LL of graphene (the sublattice-valley locking occurs because the wave functions of the 0LL in each valley, \( K \) or \( K' \), reside solely on one of the sublattices, A or B, of graphene [10]). According to the results of the atomic-scale STS maps, we can also exclude spin-polarized ferromagnetic or antiferromagnetic states, which should exhibit a hexagonal honeycomb density distribution in real space [14], and Kekulé bond states with the electron density only on C–C bonds [24]. The CDW phase is observed under two scenarios. In one scenario, the Coulomb interaction is strongly screened, and the observed splitting of the 0LL is much smaller than the on-site Coulomb energy. Then, the CDW phase should be observed in charge-neutral graphene, as reported very recently in graphene on a high dielectric constant substrate [26]. In our experiment, the observed charge gap of the 0LL agrees quite well with the on-site Coulomb energy for 8.2 T < \( B < 11.5 \, \text{T} \). Therefore, we can exclude this scenario. In the other scenario, a moiré superlattice can add sublattice symmetry-breaking in graphene to stabilize the CDW phase, as observed in graphene on hexagonal boron nitride with moiré superlattices [27]. In our device, the moiré superlattice structure generated between WSe\(_2\) and graphene may help stabilize the CDW phase at charge-neutral graphene for 8.2 T < \( B < 11.5 \, \text{T} \).

Along with the anomalous field-dependent charge gap in the half-filled 0LL for \( B > 11.5 \, \text{T} \), our detailed measurements of the ground states as a function of the magnetic field reveal phase transitions with continuously changed valley polarization and, even more unexpected, valley inversion (Figure 2). The 0LL\(^-\) (0LL\(^+\)) is mainly polarized in the A (B) sublattice at 10 T, as shown in Figure 2(a) (left panels) but becomes mainly polarized in the B (A) sublattice at 12.2 T (Figure 2(a), middle panels). By further increasing the field to 13.5 T, the 0LL\(^-\) (0LL\(^+\)) changes back to localize mainly in the A (B) sublattice (Figure 2(a), right panels). Figure 2(b) shows representative line cuts of the conductance maps measured under different magnetic fields (more experimental results in Figure S5). The changes in sublattice polarization with the magnetic field for \( B > 11.5 \, \text{T} \) indicate a field-induced valley inversion due to the sublattice-valley
locking in the 0LL of graphene. To quantitatively analyze the field-dependent sublattice (valley) polarization, we plotted $Z = \frac{I_A - I_B}{I_A + I_B}$ as a function of the magnetic field in Figure 2(c) (here $I_A$ and $I_B$ are the intensity of conductance maps at the A and B sublattices, respectively, as shown in Figure 2(b)). The magnitude of $Z$ reflects the degree of sublattice (valley) polarization [27]. For $Z=\pm 1$, we obtain a fully valley-polarized phase. In contrast, when $Z=0$, the valley is completely unpolarized, indicating that the spin is polarized [6-13]. For $8.2 \text{T} < B < 11.5 \text{T}$, we observe the valley-polarized CDW phase; however, the measured valley polarization is less than one. This discrepancy is mainly due to the overlap of the 0LL− and 0LL+, which have a line width of approximately 20-30 meV and are separated by approximately 20-30 meV. Therefore, we always observe a nonzero value on the A (B) sites in the conductance map of the 0LL+ (0LL−), even though it is predicted to be zero in the valley-polarized CDW phase. By increasing the magnetic field strength to $B > 11.5 \text{T}$, the valley polarization decreases and reverses its direction at $B = 12.2 \text{T}$. The inverse valley polarization reaches its maximum at approximately 12.8 T and decreases to zero again at approximately 13.2 T. By further increasing the magnetic field strength to $B > 13.2 \text{T}$, the valley polarization returns to its initial direction. After reaching the maximum at approximately 13.6 T, the valley polarization decreases to almost zero at $B = 14 \text{T}$. Obviously, our measurements reveal phase transitions with continuous changes in valley polarization and valley inversion for $B > 11.5 \text{T}$, which enriches the phase diagram of the QHFM beyond current theoretical predictions. The observed phase transition may affect the lattice-scale symmetry-breaking terms and result in an anomalous field-dependent charge gap in the half-filled 0LL for $B > 11.5 \text{T}$, as shown in Figure 1(d).

In the graphene/WSe2 heterostructure, the WSe2 nanostructures at the interface between graphene and WSe2 locally introduced the electrostatic potential to change the doping of graphene [29,30], allowing us to locally study the effect of doping on the splitting of the 0LL. Figure 3(a) shows a representative STM image of the graphene/WSe2 heterostructure with a nanoscale WSe2 quantum dot (QD) at the interface (Figure 3(b) for a schematic of the structure). The WSe2 QD generates a smooth electrostatic potential to bend the LLs and lift the orbital degeneracy of the LLs around the WSe2 QD, as shown in Figure 3(c). No intervalley scattering occurs in graphene around the WSe2 QD because the electrostatic potential is quite smooth. For graphene above the nanoscale WSe2 QD, there are electron whispering-gallery modes and atomic collapse states due to the confinement of the local electrostatic potential (Figure S6 for a detailed discussion) [29,30]. The bending of the LLs changes the doping of graphene and completely removes the splitting of the 0LL when the filling factor is changed from 0 to −2 (the 0LL is changed from half-full to empty), as shown in Figure 3(c) and (d). The result that the splitting of the 0LL seems to close for a large distance in Figure 3(c) is caused by the slight variation in the doping in graphene. This result further confirms the Coulomb interaction as the origin of the splitting of the half-filled 0LL. The bending of the LLs generates a quantum Hall conducting channel in graphene around the WSe2 QD. Figure 3(e) shows the STS map in graphene around the WSe2 QD at the Fermi energy ($E=0 \text{meV}$). A ring-shaped structure is observed in graphene around the WSe2 QD because of the 0LL across the Fermi level. This

Figure 2 (Color online) (a) STS maps of the 0LL− and 0LL+ states at $\nu=0$. The hexagonal honeycomb structures of graphene are overlaid onto the STS maps. At $B = 10 \text{T}$, the CDW phase is observed. At $B = 12.2 \text{T}$ and $13.5 \text{T}$, the maps reveal a phase transition with unexpected valley polarization and inversion from the CDW phase. (b) Normalized intensity of STS signals along the black lines (0LL−) and green lines (0LL+) in (a), showing intensities in the sublattice A and B. (c) Sublattice polarization versus the magnetic field.
conducting channel is approximately 5-10 nm in width and is quite robust when the bent 0LL is across the measured energy level (Figure S7). For the STS maps measured at the energies of the 0LL\(^-\) and 0LL\(^+\), we can also obtain the distribution of the two states in graphene around the WSe\(_2\) QD (Figure S7). For STS maps at energies between the (bent) LLs, we obtain a uniform insulating feature in graphene around the WSe\(_2\) QD. Our result demonstrates that the local electrostatic potential can generate a quantum Hall conducting channel inside a quantum Hall insulator, which is directly imaged using an STM. This finding indicates that the quantum interference and excitation of the quantum Hall conducting channels can be studied using STM.

3 Conclusions

In summary, we report the field-tuned continuous phase transition of different valley orderings in the QHFM of charge-neutral graphene on insulating WSe\(_2\). Using atomic resolution imaging of electronic wavefunctions, we observe the microscopic signatures of field-tuned continuous-varied valley polarization and valley inversion. Additionally, the phase transition is clearly revealed by an anomalous field-dependent energy gap in the half-filled 0LL. These results are unexpected and enrich the phase diagram of the QHFM beyond current theoretical predictions.

Supporting Information

The supporting information is available online at http://phys.scichina.com and https://link.springer.com. The supporting materials are published as submitted, without typesetting or editing. The responsibility for scientific accuracy and content remains entirely with the authors.

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1 L. Balents, C. R. Dean, D. K. Efetov, and A. F. Young, Nat. Phys. 16, 725 (2020).
2 E. Y. Andrei, and A. H. MacDonald, Nat. Mater. 19, 1265 (2020), arXiv: 2008.08129.
3 Y. N. Ren, Y. Zhang, Y. W. Liu, and L. He, Chin. Phys. B 29, 117303 (2020), arXiv: 2008.09769.
4 E. Y. Andrei, D. K. Efetov, P. Jarillo-Herrero, A. H. MacDonald, K. F. Mak, T. Senthil, E. Tutuc, A. Yazdani, and A. F. Young, Nat. Rev. Mater. 6, 201 (2021).
5 Q. Zheng, C. Y. Hao, X. F. Zhou, Y. X. Zhao, J. Q. He, and L. He, Phys. Rev. Lett. 129, 076803 (2022), arXiv:2207.12670.
6 K. Yang, S. Das Sarma, and A. H. MacDonald, Phys. Rev. B 74, 075423 (2006), arXiv: cond-mat/0605666.
7 J. Alicea, and M. P. A. Fisher, Phys. Rev. B 74, 075422 (2006), arXiv: cond-mat/0604601.
8 K. Nomura, and A. H. MacDonald, Phys. Rev. Lett. 96, 256602 (2006), arXiv: cond-mat/0604113.
9 J. Jung, and A. H. MacDonald, Phys. Rev. B 80, 235417 (2009), arXiv: 0909.1362.
10 M. Kharitonov, Phys. Rev. B 85, 155439 (2012), arXiv: 1103.6285.
11 T. Jolicoeur, and B. Pandey, Phys. Rev. B 100, 115422 (2019), arXiv:1907.04612.
12 J. Atteia, Y. Lian, and M. O. Goerbig, Phys. Rev. B 103, 054503 (2021), arXiv:2011.11830.
13 J. G. Checkelsky, L. Li, and N. P. Ong, Phys. Rev. Lett. 100, 206801 (2008), arXiv:0708.1959.
14 A. F. Young, C. R. Dean, L. Wang, H. Ren, P. Cadden-Zimansky, K. Watanabe, T. Taniguchi, J. Hone, K. L. Shepard, and P. Kim, Nat. Phys. 8, 550 (2012), arXiv: 1201.4167.
15 A. F. Young, J. D. Sanchez-Yamagishi, B. Hunt, S. H. Choi, K. Watanabe, T. Taniguchi, R. C. Ashoori, and P. Jarillo-Herrero, Nature 505, 528 (2014), arXiv: 1307.5104.
16 A. A. Zibrov, E. M. Spanton, H. Zhou, C. Kometter, T. Taniguchi, K. Watanabe, and A. F. Young, Nat. Phys. 14, 930 (2018), arXiv: 1712.01968.
17 Y. Zhao, P. Cadden-Zimansky, F. Ghahari, and P. Kim, Phys. Rev. Lett. 108, 106804 (2012), arXiv: 1201.4434.
18 A. Knothe, and T. Jolicoeur, Phys. Rev. B 92, 165110 (2015), arXiv: 1507.05866.
19 V. P. Gusynin, V. A. Miransky, S. G. Sharapov, and I. A. Shovkovy, Phys. Rev. B 77, 205409 (2008), arXiv: 0806.2136.
20 V. P. Gusynin, V. A. Miransky, S. G. Sharapov, I. A. Shovkovy, and C. M. Wyenberg, Phys. Rev. B 79, 115431 (2009), arXiv: 0801.0708.
21 P. K. Pyatkovskiy, and V. A. Miransky, Phys. Rev. B 90, 195407 (2014), arXiv: 1409.1629.
22 S. Kim, J. Schwenk, D. Walkup, Y. Zeng, F. Ghahari, S. T. Le, M. R. Slot, J. Berwanger, S. R. Blankenship, K. Watanabe, T. Taniguchi, F. J. Giessibl, N. B. Zhitenev, C. R. Dean, and J. A. Stroscio, Nat. Commun. 12, 2852 (2021), arXiv: 2006.10730.
23 K. Lai, W. Kundhikanjana, M. A. Kelly, Z. X. Shen, J. Shabani, and M. Shayegan, Phys. Rev. Lett. 107, 176809 (2011), arXiv: 1110.0067.
24 S. Y. Li, Y. Zhang, L. J. Yin, and L. He, Phys. Rev. B 100, 085437 (2019), arXiv: 1904.06902.
25 L. Veyrat, C. Déprez, A. Coissard, X. Li, F. Gay, K. Watanabe, T. Taniguchi, Z. Han, B. A. Piot, H. Sellier, and B. Sacépé, Science 367, 781 (2020), arXiv: 1907.02299.
26 A. Coissard, D. Wander, H. Vignaud, A. G. Grushin, C. Repellin, K. Watanabe, T. Taniguchi, F. Gay, C. B. Winkelmann, H. Courtois, H. Sellier, and B. Sacépé, Nature 605, 51 (2022), arXiv: 2110.02811.
27 X. Liu, G. Farahi, C. L. Chiu, Z. Papic, K. Watanabe, T. Taniguchi, M. P. Zaletel, and A. Yazdani, Science 375, 321 (2022), arXiv: 2109.11555.
28 Y. N. Ren, M. H. Zhang, C. Yan, Y. Zhang, and L. He, Sci. China-Phys. Mech. Astron. 64, 287011 (2021).
29 Q. Zheng, Y. C. Zhuang, Q. F. Sun, and L. He, Nat. Commun. 13, 1597 (2022), arXiv: 2110.06673.
30 Q. Zheng, Y. C. Zhuang, Y. N. Ren, C. Yan, Q. F. Sun, and L. He, Phys. Rev. Lett. 130, 076202 (2023), arXiv: 2206.07221.
31 Y. Jiang, X. Lai, K. Watanabe, T. Taniguchi, K. Haule, J. Mao, and E. Y. Andrei, Nature 573, 91 (2019), arXiv: 1904.10153.
32 Y. Choi, J. Kemmer, Y. Peng, A. Thomson, H. Arora, R. Polski, Y. Zhang, H. Ren, J. Alicea, G. Refael, F. von Oppen, K. Watanabe, T. Taniguchi, and S. Nadji-Perge, Nat. Phys. 15, 1174 (2019), arXiv: 1901.02997.
33 A. Kerelsky, L. J. McGilly, D. M. Kennes, L. Xian, M. Yankowitz, S. Chen, K. Watanabe, T. Taniguchi, J. Hone, C. Dean, A. Rubio, and A. N. Pasupathy, Nature 572, 95 (2019).
34 Y. Xie, B. Lian, B. Jäck, X. Liu, C. L. Chiu, K. Watanabe, T. Taniguchi, B. A. Bernevig, and A. Yazdani, Nature 572, 101 (2019), arXiv: 1906.09274.
35 S. Y. Li, Y. Zhang, Y. N. Ren, J. Liu, X. Dai, and L. He, Phys. Rev. B 102, 121406 (2020), arXiv: 1912.13133.
36 Z. Q. Fu, K. K. Bai, Y. N. Ren, J. J. Zhou, and L. He, Phys. Rev. B 101, 235310 (2020), arXiv: 2001.02493.
37 A. Laturia, M. L. Van de Put, and W. G. Vandenberghe, npj 2D Mater. Appl. 2, 6 (2018).
38 H. S. Arora, R. Polski, Y. Zhang, A. Thomson, Y. Choi, H. Kim, Z. Lin, I. Z. Wilson, X. Xu, J. H. Chu, K. Watanabe, T. Taniguchi, J. Alicea, and S. Nadji-Perge, Nature 583, 379 (2020).
39 Y. N. Ren, Q. Cheng, S. Y. Li, C. Yan, Y. W. Liu, K. Lv, M. H. Zhang, Q. F. Sun, and L. He, Phys. Rev. B 104, L161408 (2021), arXiv: 2103.16127.
40 Y. N. Ren, Q. Cheng, Q. F. Sun, and L. He, Phys. Rev. Lett. 128, 206805 (2022).