Hofstadter subband ferromagnetism and symmetry-broken Chern insulators in twisted bilayer graphene

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When the twist angle between two layers of graphene is approximately 1.1°, interlayer tunnelling and rotational misalignment conspire to create a pair of flat bands1 that are known to host various insulating, superconducting and magnetic states when they are partially filled4,5. Most work has focused on the zero-magnetic-field phase diagram, but here we show that twisted bilayer graphene in a finite magnetic field hosts a cascade of ferromagnetic Chern insulators with Chern number $|C| = 1, 2$ and 3. The emergence of the Chern insulators is driven by the interplay of the moiré superlattice with the magnetic field, which endows the flat bands with a structure of topologically non-trivial subbands characteristic of the Hofstadter butterfly6–9. The new phases can be accounted for in a Stoner picture10, in contrast to conventional quantum Hall ferromagnets, electrons polarize into between one and four copies of a single Hofstadter subband11,12. Distinct from other moiré heterostructures13–15, Coulomb interactions dominate in twisted bilayer graphene, as manifested by the appearance of Chern insulating states with spontaneously broken superlattice symmetry at half filling of a $C = -2$ subband16,17. Our experiments show that twisted bilayer graphene is an ideal system in which to explore the strong-interaction limit within partially filled Hofstadter bands.

The energy spectrum of a two-dimensional electron subject to external magnetic and periodic potential is described by a fractal structure known as the Hofstadter butterfly9–11. For a system with $n$ electronic bands at zero magnetic field, the Hofstadter spectrum at a magnetic flux per unit cell of $\Phi_0/\Phi_0 = p/q$ hosts $nq$ ‘subbands’ (here $\Phi_0 = h/e$ is the magnetic flux quantum and $p$ and $q$ are integers with greatest common divisor 1). Remarkably, these subbands are each characterized by a non-zero, integer-valued topological index, known as a Chern number, $C$ which describes their contribution to the quantized Hall conductivity when an integer number of subbands are filled16. Owing to their high sample quality and the large area of the superlattice, moiré van der Waals heterostructures have provided a versatile materials system for exploring the physics of the Hofstadter subbands13–15, including the observation of correlation-driven Chern insulators at fractional filling of a Hofstadter subband that either spontaneously break the superlattice symmetry15,18 or fractionalize electrons into anyons15. In typical moiré heterostructures of graphene and hexagonal boron nitride (hBN), however, accessing the Hofstadter regime experimentally requires strong magnetic fields. This is tied to the fact that the superlattice potential induced by the moiré pattern is weak, and provides only a small perturbation to the conventional Landau levels. In this limit, Hofstadter subbands form within a single Landau level, and are restricted to an energy bandwidth per Landau level of $W \propto \exp(-\frac{C}{\theta_0})$ (ref. 18). Achieving separations between subbands that exceed the energy scales related to disorder or temperature thus requires a substantial fraction of the magnetic flux quantum to be threaded through each superlattice unit cell.

Electrons in the twisted bilayer graphene (tBLG) flat bands are strongly localized and should consequently realize a strong-lattice limit of the Hofstadter butterfly in a finite magnetic field. In this limit, topologically non-trivial subbands may be separated by a substantial fraction of the total bandwidth even for $\Phi_0 < \Phi_0$ (refs. 11,12,19). However, discussions in the literature of the magneto-transport properties of tBLG have focused on relating the observed magnetoresistance features to the enigmatic magnetic field ($B = 0$) phase diagram, without regard for the interplay of the magnetic flux with the periodic potential. Here, we show that magnetic fields can indeed drive new symmetry-breaking transitions, producing magnetoresistance features at low magnetic fields that are not correlated with the realized $B \approx 0$ ground states. Figure 1a shows longitudinal resistivity ($\rho_{xx}$) data from a nominally hBN-unaligned twist angle $\theta = 1.12°$ tBLG device (Extended Data Fig. 1) as a function of $B$. We present data at 4 K to suppress all but the most energetically robust features; lower-temperature $\rho_{xx}$ data and Hall conductivity ($\sigma_{xy}$) data are available in Extended Data Figs. 2 and 3. At $B = 0$ T, our device shows familiar features of flat-band tBLG: correlated insulators or magnetoresistance peaks at $\nu \approx \pm 2, \pm 3$ and 0 and a robust superconducting state for $\nu \approx -2$ (Fig. 1c and Extended Data Fig. 4). Here $\nu$ indicates the number of electrons per superlattice unit cell. However, at $B \approx 5$ T (for $\nu < 0$) and $B \approx 2$ T (for $\nu > 0$) this behaviour abruptly changes, giving way to a series of Chern insulators, indicated in blue in Fig. 1b. They are characterized by $\rho_{xx}$ minima and integer quantized Hall conductivity that increases in magnitude as the band is filled.

Within the Hofstadter picture, gapped states follow linear trajectories according to a Diophantine equation, $\nu = t n_k + s$ ($s, t \in Z$) in the space of $\nu$ and $n_k \equiv \Phi_0/\Phi_0$ (ref. 20), the magnetic flux per moiré unit cell. Here $s$ is the Bloch band filling index, which encodes the number of electrons bound to each lattice unit cell, while $t$ is the total Chern number associated with a given gap, and consequently the quantized Hall conductivity21,22. Starting from $\nu = -4$ and moving into the flat band, at $n_k = 0.3$ the observed sequence follows ($t, s$)

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Fig. 1 | Hofstadter subband ferromagnetism. a, $\rho_{xx}$ as a function of $\nu$ and $B$ at 4 K. b, Schematic of the ferromagnetic Chern insulator structure observed in a, showing gapped states that follow linear trajectories $\nu = n_0 + s$. We estimate $n_0 = 1$ for $B = 30.1 \pm 0.1 T$. Two trajectory classes are distinguished by colour corresponding to $s = 0, 1 \in Z$ integer quantum Hall states (IQH) and $s, t \in Z, s \neq 0$ fermionic Chern insulators (CI). c, $\rho_{xx}$ (black) and $\rho_{xy}$ (red) as functions of $\nu$ at 0.6 K at 0 T ($\rho_{xx}$) and 0.6 T ($\rho_{xy}$). Schematic images show the pattern of flavour filling consistent with observed low-magnetic-field quantum oscillations (Supplementary Information). d, $\rho_{xx}$ at $T = 4 K$ (black) and $\sigma_{xy}$ at 4 K (red) and 10 mK (blue) as a function of $\nu$ at $B = 10 T$. The dashed lines show integer multiples of $e^2/h$. Schematic images show the sequence of filling of the Hofstadter subbands. e, Calculated DOS of the valence and conduction flat bands at zero magnetic field. a.u., arbitrary units. f, Calculated Hofstadter spectra of tBLG. The blue and yellow regions correspond to Chern bands with total $C = -1$ and +2, respectively. (See Supplementary Information for details.) g, Calculated DOS at $n_e = 2/5$.

$=(-1,-3), (-2,-2), (-3,-1)$ and $(-4,0)$. This pattern is consistent with sequentially filling four bands with $C = -1$, each of which binds a single electron per moiré unit cell.

The observed transition between low- and high-magnetic-field regimes can be understood from the evolution of the tBLG flat bands in a magnetic field. At $B = 0$, tBLG hosts two low-energy bands (for a given spin and valley flavour) connected by two Dirac points that remain gapless absent symmetry breaking by substrate potentials21–24. A key feature of the $B = 0$ band structure is the triangular form of the density of states (DOS), shown in Fig. 1c. It was recently argued25 that exchange interactions within such a band lead to a peculiar form of Stoner ferromagnetism in which exchange interactions favour collective states that are symmetric between two, three or four isospin flavours. In our device, the low-$B$ quantum oscillations (Extended Data Fig. 5) are qualitatively consistent with this mechanism, showing signatures of Dirac fermions with fully broken flavour symmetry for $\nu < -3$, two-fold symmetry for $-3 < \nu < -2$ and four-fold symmetry for $-2 < \nu < 0$ (insets to Fig. 1c and Extended Data Fig. 5). Notably, in our device, and indeed in the majority of published transport data in tBLG devices, symmetry breaking is absent at $\nu = \pm 1$, as evidenced by the absence of low-$B$ quantum oscillations near these fillings. This effect is not explained within the model of ref. 25, but may arise due to the reconstruction of the low-energy bands that is expected to occur at partial fillings due to the Coulomb interactions26.

An out-of-plane magnetic field substantially reconstructs the low-energy bands. Figure 1f shows the evolution of electronic structure as a function of magnetic field, while Fig. 1g shows the calculated DOS at $n_e = 2/5$, corresponding to $B = 12 T$ for our device. At low magnetic fields, the magnetic field reconstructs the two flat bands into three groups of subbands (for each spin and valley), having net Chern numbers $-1$, $+2$ and $-1$ respectively. This structure is generic to all tBLGs, with the $C = 2$ subband originating in the zero-energy Landau level in the two constituent graphene monolayers and the $C = -1$ bands ensuring that the combined bands maintain a net Chern number of zero. Even in moderate fields the DOS is characterized by three effective subbands, in contrast to the two bands at $B = 0$. The cascade of Chern insulators is well explained by the sequential filling of four spin- and valley-projected copies of the lower $C = -1$ subbands, as illustrated in Fig. 1d, and emerges from a simplified analysis of a Stoner model that accounts for the magnetic-field-induced change in bandwidth (Supplementary Fig. 3). This picture is mirrored for electron doping, and is also consistent with the Landau fan originating from the charge neutrality point, which shows the strongest features at $(t, s) = (\pm 4, 0), (\pm 2, 0)$ and $(0, 0)$, consistent with successive fillings of the central $C = 2$ subband.

The most striking features of the ferromagnetic Chern insulator cascade are the states at $(t, s) = (\pm 3, \pm 1)$ (Fig. 2 and Extended Data Fig. 6), which are not correlated with a set of low-magnetic-field quantum oscillations (Extended Data Fig. 5). Indeed, low-temperature measurements show that these states emerge without warning in the midst of well formed four-component quantum oscillations originating from the charge neutrality point, consistent
Fig. 2 | Magnetic-field-driven symmetry-breaking transition to a C = −3 Chern insulator. a–b, \( \rho_{xx} \) (a) and \( \sigma_{xx} \) (b) as functions of \( \nu \) and \( B \) between 4 and 6.5 T, measured at a nominal temperature of 10 mK. c, Line cuts of \( \rho_{xx} \) and \( \sigma_{xx} \) along yellow lines in a and b. Dashed lines correspond to \( -3e^2/h \) and \(-Bx^2/h\). Schematics depict inferred band fillings for \((\nu, s) = (-8, 0) \) and \((-3, -1)\). At \((-3, -1)\), three \( C = -1 \) subbands are filled and one is empty. At \((-8, 0)\) all four flavours are equally occupied, filling four copies of a \( C = -2 \) band including the entire \( C = -1 \) band except for the highest-energy state associated with a single, strain-split Landau level\(^\text{12}\).

with previous reports of orbital magnetic states at finite \( B \) (refs. 5,27). The sudden phase transition (Fig. 2a,b) and sudden appearance of an activation gap (Extended Data Fig. 7) can be understood within the context of Stoner ferromagnetism, which accounts for the Chern character of the \( C = -1 \) Hofstadter subband. Within this picture, symmetry breaking occurs when DOS > \( 1/U \), where \( U \) is the exchange interaction strength. As the magnetic field increases, both the bandwidth of and total number of states within the \( C = -1 \) subband decrease; numerically, however, this leads to an increase in average DOS (Supplementary Fig. 2), favouring symmetry breaking above some critical \( B \). Here, we did not observe any hysteric behaviour in this transition. The nature of this first-order-like transition is illustrated in Fig. 2c, which shows a line cut connecting the \((-3, -1)\) and \((-8, 0)\) states. At \((-3, -1)\) three \( C = -1 \) subbands are completely filled and one empty, while at \((-8, 0)\) all four flavours are equally occupied, filling four copies of a \( C = -2 \) band that includes the entire \( C = -1 \) band except for the highest-energy state associated with a single, strain-split Landau level\(^\text{12}\).

In one view, the ferromagnetic Chern insulator states can be seen as competing ground states in tBLG at integer band filling and \( B = 0 \) (refs. 30,31). It is known that small perturbations such as an aligned hBN substrate can favour Chern insulators at \( B = 0 \), for instance at \( \nu = 3 \) (ref. 3). Even absent this splitting, however, a retrospective review shows that nearly all published data containing magneto-transport data\(^\text{3,4,7,31}\), as well as recent studies\(^\text{32}\), show a feature probably associated with the transition we report here to one or more of the Chern insulator states we observe. These transitions can be generically understood as reflecting the finite orbital moment of the Chern insulator states, which lowers their energy relative to competing \( C = 0 \) insulators or semimetallic ground states in a finite magnetic field. However, we prefer to analyse the observed ferromagnetic Chern insulator states as generalizations of quantum Hall ferromagnets in which the Landau levels are replaced by \( C = -1 \) subbands of the tBLG Hofstadter spectrum. We note that, just as in graphene Landau levels\(^\text{33}\), each Chern ferromagnet corresponds to polarization within the combined space of spins and valleys with the precise isospin ordering set by the interplay of higher-order effects including the Zeeman effect and the anisotropy of the Coulomb

Fig. 3 | SBCIs. a–b, \( \rho_{xx} \) (a) and \( \sigma_{xx} \) (b) as a function of \( \nu \) and \( B \) in a range of 6–12 T at a nominal temperature of 10 mK. c, Schematic of observed Chern insulator states in a and b. Red lines correspond to SBCIs with \((t, s) = (-2, -3/2) \) and \((-3, -1/2)\). d, Calculated Hofstadter energy spectrum in the regime of e, focusing on the low-energy \( C = -1 \) band. The two primary subbands are colour coded, and the equivalent extent in density is shown in e for two copies of the band. e, Line cuts of \( \rho_{xx} \) and \( \sigma_{xx} \) at 11.2 T with several Chern insulator states labelled.
interactions themselves. Interestingly, these anisotropies may differ within the Chern bands as compared with both Landau levels and the $B=0$ flat bands, possibly favouring different broken symmetries at integer fillings.

This approach moreover has the advantage of providing a quantitative description of the electronic substructure of the $C=x=1$ band at high magnetic fields. Figure 3a,b displays measured $\rho_{xx}$ and $\sigma_{xy}$, respectively, from 6 to 12 T in a density region corresponding to occupying the third and fourth copies of the lower $C=-1$ bands. The low disorder in our graphite gated sample allows us to resolve a number of smaller energy gaps characterized within the $C=x=1$ band. First, we resolve the next level of subband structure within flat-band systems the magnetic-field-independent moiré wavefunction $B$ which vanishes at low fields. In contrast, in tBLG and other moiré structures the Coulomb interaction scale is never realized in tBLG. Specifically, within the Hofstadter subbands of a related to the strong-lattice limit of the Hofstadter band structure appear only at magnetic fields in excess of 16 T, more than twice the low disorder in our graphite gated sample allows us to resolve occupying the third and fourth copies of the lower $C=\pm 1$ band described above. In addition, we also observe robust states showing $\langle t,s\rangle=\langle -3,-1/2 \rangle$ and $\langle -2,-3/2 \rangle$ at half filling of the $C=x=2$ band (highlighted in red in Fig. 3c). These states show quantized $\sigma_{xy}=e^{2}\hbar$ as expected from the measured slope $t$, and an activation gap of around 1 K (Fig. 3e and Extended Data Figs. 8 and 9). Similar states are observed at positive fillings as well (Extended Data Fig. 8).

The quantum number $s$ encodes the total number of electrons bound to a given lattice site at fixed magnetic flux, and is an integer for all states that can be described within a non-interacting Hofstadter butterfly. The simplest mechanism that allows this number to be fractional is for the superlattice symmetry to break, and we associated the observation of states with half-integer $s$ with the formation of symmetry-broken Chern insulators (SCIs), in which electronic interactions within a $C=x=2$ Hofstadter subband drive spontaneous doubling of the unit cell. While such states were previously observed in hBN-aligned graphene, they typically appear only at magnetic fields in excess of 16 T, more than twice as high as in tBLG. The relative strength of the SCI states can be related to the strong-lattice limit of the Hofstadter band structure realized in tBLG. Specifically, within the Hofstadter subbands of a weak-superlattice system the Coulomb interaction scale is never stronger than $e^{2}/\epsilon_{G}$ (here $\epsilon_{G}=\sqrt{\hbar/eB}$ is the magnetic length), which vanishes at low $B$. In contrast, in tBLG and other moiré flat-band systems the magnetic-field-independent moiré wave-length sets the scale for interactions. Interactions can drive the formation of correlated states at partial filling of Hofstadter subbands in the low-field limit, just as they drive correlated states at $B=0$. Our experiment suggests that other exotic ground states at fractional fillings of Chern bands may be accessible in moiré flat-band systems, potentially even at $B=0$ in the presence of appropriate time-reversal symmetry-breaking order.

Online content
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Methods
Device fabrication. The tBLG device used in this study was fabricated using the ‘cut-and-stack’ technique described in ref. 34; and indeed the device is the same as device number 5 in that paper. Before stacking, we first cut the graphene into two pieces using atomic force microscopy to prevent unintentional strain in tearing graphene. We used a poly(bisphenol A carbonate) (PC)/polymethylmethacrylate stamp mounted on a glass slide to pick up a 40-nm-thick hBN flake at 90–110°C, and carefully pick up the first half of a precut graphene piece, rotate the sample stage and pick up the second half of the graphene at 25°C using the hBN flake. We rotated the graphene pieces manually by a twist angle of about 1.2–1.3° before the second pickup. Finally, the two-layer stack (hBN–tBLG) is transferred onto another stack (40-nm-thick hBN–graphite) containing the bottom gate and gate dielectric, which is prepared in advance by the same dry transfer process and cleaned with a typical solvent wash using chloroform, acetone, methanol and isopropyl alcohol followed by vacuum annealing (400°C for 8 h) to remove the residue of the PC film on the hBN surface. Electrical connections to the tBLG were made by CHF3/O3 etching and deposition of Cr/Pd/Au (2/15/180 nm) metal edge contacts35.

Transport measurements. The insulating nature of the ε = 2, 3 and 4 states appears to be correlated with poor contact resistances $R > 1 \mathrm{M} \Omega$ for $\nu \geq 3$. We thus focus the bulk of our analysis on negative filling. All transport measurements in this study were carried out in a top-loading dilution refrigerator (Bluefors LD400) with a nominal base temperature of 10 mK, which is equipped with a 14 T superconducting magnet and heavy radiofrequency and audio-frequency filtering with a cutoff frequency of ~10 kHz. The temperature-dependent measurements were made by controlling the temperature using a heater mounted on a mixing-chamber plate. Standard low-frequency lock-in techniques with Stanford Research SR860 amplifiers were used to measure the resistance with an excitation current of $2 \times 10^{-8} \text{A}$ at a frequency of 17.777 Hz. $\theta$ is determined from the values of charge carrier density at which the insulating states at $n_{\text{LL}}=\varepsilon$ are observed, following $n_{\text{LL}}=\varepsilon^{\pm}d/\sqrt{3}a^2$, where $d=0.246$ nm is the lattice constant of graphene.

Low-field quantum oscillation analysis. The low charge and twist-angle disorder of our device allows us to resolve quantum oscillations down to very low magnetic fields. Extended Data Fig. 5 shows quantum oscillation data for $-4 < \nu < 0$. At a magnetic field of 0.6 T, quantum oscillations are not visible except originating from the charge neutrality point. Notably, the strongest features are at Landau level filling $n_{\text{LL}}=4$, ±12, ±20 (Extended Data Fig. 5c), precisely as expected from the non-interacting band structure of tBLG35. This sequence can be understood as twice that of monolayer graphene, and indeed arises from the two gapless Dirac cones of the constituent graphene monolayers.

This is in marked contrast to published quantum oscillations in other twisted bilayers, which uniformly show a sequence of ±4, ±8, ±12 and so on. However, we see the emergence of the more typical sequence at higher magnetic fields, for example in the data taken at $B=2$ T shown in Extended Data Fig. 5c. By this magnetic field, we also see quantum oscillations near $\nu=2$ and $\nu=1$, with sequences of $-4$, $-2$, $-6$, ... and $-1$, $-2$, $-3$, $-4$, ..., respectively. The transition in behaviour may arise from a single-particle effect such as strain35, which may well be ubiquitous in graphene devices at some level. However, conventional quantum Hall ferromagnetism may also be at play, as evidenced by the emergence of additional quantum oscillation minima at $n_{\text{LL}}=1, 2, 3$ and 6.

The data can be elegantly explained under an ansatz consistent with that taken in ref. 34, with the modification that no three-fold symmetry state exists.

Data availability
Source data are provided with this paper. All other data that support the plots within this paper and other findings of this study are available from the corresponding author upon reasonable request.

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Author contributions
Y.S. and J.G. fabricated tBLG devices. Y.S. performed the measurements and analysed the data. L.R. and D.A.A. performed the theoretical calculations. Y.S. and A.F.Y. wrote the paper with input from L.R. and D.A.A. T.T. and K.W. grew the hBN crystals.

Competing interests
The authors declare no competing interests.

Additional information
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Extended Data Fig. 1 | Device characterization. a, Optical microscope images of a device used in this study. Scale bar equals 5μm. b, Two terminal conductance $G_2$ across multiple contacts at 4 K showing high degree of uniformity. ‘All’ denotes all other contacts.
Extended Data Fig. 2 | Landau fan diagram up to 13.5 T. a, $\rho_{xx}$ as a function of $\nu$ and $B$ at a nominal temperature of 10 mK. b, Schematic of observed Chern insulator structure, showing gapped states that follow linear trajectories parameterized by $\nu = n_s + s$. $n_s$ and $\nu$ are the magnetic flux quanta and number of electrons per moiré unit cell, respectively. Three trajectory classes are distinguished by color: Integer quantum Hall (IQH) (gray; $s=0, t \in \mathbb{Z}$), Chern insulators (CI) (blue; $s, t \in \mathbb{Z}$, $s \neq 0$) and symmetrybroken Chern insulators (SBCI) (magenta; fractionals, $t \in \mathbb{Z}$).
Extended Data Fig. 3 | Landau fan diagram of Hall conductivity. a, b, $\sigma_{yx}$ as a function of $\nu$ and $B$ at 4 K (a) and 10 mK (b).
Extended Data Fig. 4 | Superconductivity near $\nu = -2$. a, $\rho_{xx}$ as a function of $\nu$ and $T$ around $\nu = -2$. b, Line cut of $\rho_{xx}$ as a function of $T$ at $\nu = -2.46$. c, $\frac{dV_{xx}}{dl}$ as a function of magnetic field at a nominal temperature of 10 mK and $\nu = -2$. Both current and magnetic fields suppress the low resistance state with Fraunhofer like oscillation.
Extended Data Fig. 5 | Landau fan diagram at low magnetic fields at 0.4 K. a, $\rho_{xx}$ as a function of $\nu$ and $B$. b, Schematic observed Chern insulator structure based on a. c, Line cuts of $\rho_{xx}$ at 0.6 T in a. d, Line cuts of $\rho_{xx}$ at 2.0 T in a.
Extended Data Fig. 6 | \( C = +3 \) Chern insulator from \( \nu = 2 \).

a, \( \rho_{xx}(a) \) and \( \sigma_{yx}(b) \) as a function of \( \nu \) and \( B \) in negative fillings at 10 mK.

c, \( \rho_{xx}(c) \) and \( \sigma_{yx}(d) \) as a function of \( \nu \) in negative fillings between 0.3 and 4.7 K.

e, Line cuts of \( \rho_{xx} \) (black) and \( \sigma_{yx} \) (red) as a function of \( B \) along with yellow lines in a and b. Dashed line corresponds to \( 3e^2/h \).

f, The thermal activation gap of \( C = +3 \) Chern insulator as a function of \( B \). The values are calculated from the fits to an Arrhenius law, \( \rho_{xx} \sim \exp(-\Delta/2T) \). Error bars in the gaps represent the uncertainty arising from determining the linear (thermally activated) regime for the fit.
Extended Data Fig. 7 | Temperate dependence of a $C=-3$ Chern insulator from $\nu = -1$. a, b, $\rho_{xx}(a)$ and $\sigma_{yx}(b)$ as a function of $\nu$ in negative fillings between 0.3 and 4.7 K. c, The thermal activation gap of $C=-3$ Chern insulator as a function of $B$. The values are calculated from the fits to an Arrhenius law, $\rho_{xx} \sim \exp(-\Delta/2T)$. Error bars in the gaps represent the uncertainty arising from determining the linear (thermally activated) regime for the fit.
Extended Data Fig. 8 | Symmetry broken Chern insulators from \((t,s) = (+3,+1/2)\) and \((+2,+3/2)\). a, \(\rho_{xx}\) as a function of \(\nu\) and \(B\) in a range of 6-12 T. b, \(\sigma_{yx}\) as a function of \(\nu\) and \(B\) in a range of 6-12 T. c, Schematic observed Chern insulators corresponding to the observations in a parameterized by \(\nu = tn_\phi + s\). Red lines correspond to SBCI state with \((t,s) = (+3,+1/2)\) and \((+2,+3/2)\). d, Line cuts of \(\rho_{xx}\) and \(\sigma_{yx}\) around 10 T at positive fillings.
Extended Data Fig. 9 | Thermal activation gaps $\Delta$ of SBCIs. (a, b) of SBCI with $(t,s) = (-3,-1/2)$ and $(-2,-3/2)$ as a function of $B$. (b) $\Delta$ of SBCI with $(t,s) = (3,1/2)$ and $(2,3/2)$ as a function of $B$. 