Charge carriers in magic-angle graphene come in eight flavours described by a combination of their spin, valley and sublattice polarizations. When inversion and time-reversal symmetries are broken, this ‘flavour’ degeneracy can be lifted, and their corresponding bands can be sequentially filled. Due to their non-trivial band topology and Berry curvature, each band is classified by a topological Chern number \( C \), leading to quantum anomalous Hall and Chern insulator states. Using a scanning superconducting quantum interference device on a tip, we image the nanoscale equilibrium orbital magnetism induced by the Berry curvature, the polarity of which is governed by \( C \), and detect its two constituent components associated with the drift and self-rotation of the electronic wavepackets. At integer filling \( v = 1 \), we observe a zero-field Chern insulator, which—rather than being described by a global topologically invariant \( C \)—forms a mosaic of microscopic patches of \( C = -1, 0 \) or \( 1 \). On further filling, we find a first-order phase transition due to the recondensation of electrons from valley \( K \) to \( K' \), leading to irreversible flips of the local Chern number and magnetization, as well as to the formation of valley domain walls, giving rise to hysteretic anomalous Hall resistance.

Magic-angle twisted bilayer graphene (MATBG) exhibits a unique combination of strong electron correlations with band topology, giving rise to a wealth of emergent phenomena. Bringing the full potential of these traits to light entails two elements. The first is gapping the Dirac cones that can be attained by the application of magnetic fields or by structural modifications, such as alignment with hexagonal boron nitride (hBN). With a gapped dispersion, the singular Berry curvature at the Dirac point becomes extended over the band, giving rise to energy-dependent orbital magnetism; furthermore, the band topology is profoundly expressed in terms of non-trivial Chern gaps. The second is lifting the time-reversal symmetry; otherwise, bands related by time reversal cancel the effects of one another. In MATBG, the unique combination of gapped Dirac cones and spontaneous degeneracy lifting due to strong interactions leads to Chern insulator phases and to orbital ferromagnetic states characterized by the quantum anomalous Hall effect at full filling of a Chern band.

The Chern insulators in MATBG have been revealed by transport, scanning tunneling microscopy, and compressibility measurements. These experiments provide unambiguous evidence of gapped degeneracy-lifted bands at integer fillings that are characterized by non-trivial Chern numbers \( C \in \mathbb{Z} \), defining their topological state. Recent theoretical works, however, have proposed a new perspective on this fundamental concept: rather than being a fixed global property, a spatially varying \( C \) may occur due to position-dependent substrate potential. This additional degree of freedom opens a new platform for the construction and manipulation of previously unconsidered topological states and gapless modes. In this work, we provide the first experimental test of this hypothesis.

The Berry-curvature-induced orbital magnetism, even though extensively studied theoretically and being the main source of anomalous Hall resistivity, has evaded direct experimental observation. Moreover, the anomalous Hall effect, indicative of magnetism, has been reported in only a few MATBG samples. Recently, the first direct measurement of the local magnetization hysteresis in the gapped quantum anomalous Hall effect state in MATBG aligned to hBN was reported. On sweeping the magnetic field at \( v = 3 \), it showed a hysteretic magnetization reversal of about 3 \( \mu_0 \) per charge carrier, providing strong evidence for the orbital nature of magnetization. Here we image the continuous evolution of the equilibrium Berry-curvature-induced orbital magnetism as a function of filling factor, through both metallic and gapped states, in a non-aligned MATBG sample.

Transport measurements

The hBN-encapsulated MATBG sample with a twist angle of \( \theta \approx 1.08^\circ \) was fabricated using the cut-and-stack technique (Methods). Transport measurements of \( R_{xx} \) and \( R_{xy} \) (Fig. 1a,b) were performed in a Hall bar geometry (Fig. 2b) at temperature \( T = 300 \text{ mK} \) as a function of applied perpendicular magnetic field \( B \) and filling factor \( \nu = n/n_s = 4C_{bg}V_{bg}^d/n_s \), where \( V_{bg}^d \) is the d.c. voltage applied to the graphite back gate, \( C_{bg} \) is the back-gate capacitance, \( n \) is the carrier density and \( n_s \) corresponds to four electrons per moiré cell. The characteristic transport features of high-quality MATBG are evident including superconductivity at \( \nu = -2 - \delta \), correlated and high-field Chern insulating states, and Landau fans.
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to the graphite back gate (Fig. 2a), we set the filling factor is reversible, whereas close to
from performed minor hysteresis loops by sweeping
of magnetic field $B_c$ at $B_c$, Line cuts along the red and blue arrows in c at $B_c = 47$ mT. e, $R_{yx}$ displaying hysteresis on sweeping $B_{x}$ up (left) and down (right). f, Line cuts along the red and blue arrows in e at $v = 0.95$. g, Minor hysteresis loops of $R_{yx}$ on sweeping $v$ from 0.6 to $v_{max}$ (red) and back (blue) at $B_c = 46$ mT as $v_{max}$ increases from 0.83 to 1.30 (Supplementary Video 1). The consecutive curves are offset by $-2 \Omega$ for clarity.

Fig. 1 | Transport characteristics of MATBG. a,b. Measured $R_{xx}$ (a) and $R_{yx}$ (b) at $T = 300$ mK versus the carrier concentration ($n_v$) and applied out-of-plane magnetic field $B_{z}$, displaying superconductivity, correlated insulating states and Landau fans. c, $R_{xx}$ in the vicinity of $v = 1$ at low $B_{c}$ (black rectangle in b), displaying pronounced hysteresis on sweeping $v$ up (left) and down (right). d, Line cuts along the red and blue arrows in c at $B_c = 47$ mT. e, $R_{yx}$ displaying hysteresis on sweeping $B_{x}$ up (left) and down (right). f, Line cuts along the red and blue arrows in e at $v = 0.95$. g, Minor hysteresis loops of $R_{yx}$ on sweeping $v$ from 0.6 to $v_{max}$ (red) and back (blue) at $B_c = 46$ mT as $v_{max}$ increases from 0.83 to 1.30 (Supplementary Video 1). The consecutive curves are offset by $-2 \Omega$ for clarity.

(Extended Data Fig. 1). Notably, near $v = 1$ and at low $B_z$, non-trivial $R_{yx}$ behaviour is observed (Fig. 1c–g), indicating the presence of orbital magnetism and non-quantized anomalous Hall effect.9,10 We resolve a pronounced hysteresis in $R_{yx}$ displaying Barkhausen jumps on sweeping $B_{z}$ (Fig. 1c,f) as well as $v$ (Fig. 1c,d). Additionally, we performed minor hysteresis loops by sweeping $v$ back and forth from $v = 0.55$ to $v_{max}$ and gradually increasing $v_{max}$ (Fig. 1g and Supplementary Video 1). For $v_{max} \leq 1$, $R_{yx}$ is reversible, whereas at higher filling, hysteresis begins and grows rapidly. Notably, the full size of the hysteresis loop is attained only when $v_{max}$ exceeds 1.25, substantially larger than $v = 1.00$.

Magnetic imaging
To study the local magnetism, we utilize a scanning superconducting quantum interference device (SQUID) fabricated on the apex of a sharp pipette (SQUID on tip (SOT)). An indium SOT is with a diameter of $d = 180$ nm and field sensitivity of 10 nT Hz$^{1/2}$ is scanned at a height of $h \approx 160$ nm above the sample surface at $T = 300$ mK in the presence of a small applied $B_{c} \approx 50$ mT (Methods). By applying $V_{bg}$ to the graphite back gate (Fig. 2a), we set the filling factor near $v = 1$ and add small $V_{bg} = 20$ mV (peak to peak, square wave) with a frequency of $f \approx 6$ kHz, which modulates the carrier density by $n^{\pm}$ corresponding to filling-factor modulation of $\nu^{\pm} = 0.083$ and image the resulting $B_{z}^{\pm}(x, y) = \nu^{\pm}(dB_{z}/d\nu) = \nu^{\pm}(dB_{z}/d\nu)$ across the sample. This signal reflects the change in local magnetization due to a small change in the carrier density in the MATBG. We begin with line scans along the black dotted line (Fig. 2b). In Fig. 2f, the system is initialized at $v = 0$ and then $B_{z}^{\pm}(x, \nu)_{\uparrow}$ is measured close to $v = 1$ as $v$ is increased, $v_{\uparrow}$. A pronounced $B_{z}^{\pm}(x, \nu)_{\uparrow}$ is observed around $v = 1$, extending well below and above $v = 1$. Figure 2g shows $B_{z}^{\pm}(x, \nu)_{\uparrow}$ attained on initialization at $v = 2$ as $v$ is decreased, $v_{\downarrow}$. By subtracting the data in Fig. 2f,g, we plot (Fig. 2h) $\Delta B_{z}^{\pm}(x, \nu) = (B_{z}^{\pm}(x, \nu_{\uparrow}) - B_{z}^{\pm}(x, \nu_{\downarrow})/2$, revealing that away from $v = 1, B_{z}^{\pm}(x, \nu)_{\downarrow}$ is reversible, whereas close to $v = 1$, a substantial fraction of the line scan shows magnetization hysteresis.

Next, we perform area scans (Fig. 2b, orange rectangle), following the same $v_{\uparrow}$ and $v_{\downarrow}$ procedures and focusing closer on the vicinity of $v = 1$. The resulting $B_{z}^{\pm}(x, \nu_{\uparrow})$ and $B_{z}^{\pm}(x, \nu_{\downarrow})$ at $v = 0.966$
in Fig. 2c,d reveal a complex pattern of positive and negative $B^\infty_{x,y}$. Figure 2e shows $\Delta B^\infty_{x,y}$ that reveals the presence of magnetization hysteresis in substantial parts of the sample, whereas other areas are completely reversible at this $v$. Supplementary Video 2 presents an intricate evolution of $B^\infty_{x,y}$ on sweeping $v_1$ and $v_1$ in $v$ ranging from 0.737 to 1.174.

By performing a numerical inversion of $B^\infty_{x,y}$, we reconstruct a map of local differential magnetization $m_\nu(x,y) = dM_\nu(x,y)/dz$ (in units of $\mu_B$ per electron or equivalently $\mu_B$ per state) (Methods). Figure 3a shows that at low filling ($v=0.737$), isolated patches of positive (paramagnetic-like) $m_\nu(x,y)$ are present along with areas of negative (diamagnetic-like) response. Equivalently, $B^\infty_{x,y}$ can be inverted into current density $J^\nu(x,y)$ (in units of $\mu A \ mu m^{-1}$) (Fig. 3b). Note that magnetization $M$ in two dimensions (magnetic moment per unit area) is given in units of current, which—for the case of a uniformly magnetized domain—describes the persistent current that circulates along the edges of the domain. Figure 3b shows these sharp current channels that circulate along the domain edges.

**Tomography**

On increasing $v$, $m_\nu(x,y,v)$ evolves into a complex pattern of paramagnetic-like and diamagnetic-like domains (Fig. 3d) for $v=0.966$. The corresponding $m_\nu(x,y,v)$ and $\Delta m_\nu(x,y,v)$ are presented in Fig. 3c,d. To gain an insight into the underlying mechanism, we construct a tomographic representation of the full dataset of $m_\nu(x,y,v)$ (Fig. 3e), which shows the local $m_\nu$ evolution as a function of $v$ (Supplementary Videos 3 and 4 show the full evolutions of $m_\nu(x,y,v)$, $\Delta m_\nu(x,y,v)$ and $J^\nu(x,y,v)$ with $v$).

The tomographic dataset reveals a number of key observations. (1) The measured local differential magnetization $m_\nu$ is very large, reaching values of $25\mu_B$ per electron, which is consistent with theoretical predictions and our calculations (Methods), providing strong evidence for the orbital origin of magnetization. (2) The observation of orbital magnetism establishes valley degeneracy lifting and time-reversal symmetry breaking. The valley polarization commences well below $v=1$. (3) The pronounced variations in equilibrium local magnetization point out the local variations in band structure and Berry curvature. (4) Some regions of the sample...
show reversible magnetization, whereas others are hysteretic. (5) The hysteretic and non-hysteretic regions have comparable $m_{SR}$, suggesting that they are controlled by the same mechanism. (6) In Fig. 3, $m_{SR}(x, y)$ displays the coexistence of paramagnetic-like (red) and diamagnetic-like (blue) patches. Since paramagnetic dipole moment $m$ has favourable potential energy $U_{m} = -m \cdot B$, the presence of diamagnetic patches is surprising. (7) In conventional ferromagnets, unfavourable magnetic domains are present that minimize the $B^2$ energy and their size shrinks in a jump-wise manner on increasing $B_{0}$, giving rise to Barkhausen noise. One would, therefore, expect to see blue patches flipping abruptly to red. Remarkably, $\Delta m_{SR}(x, y)$ (Fig. 3f) shows that hysteretic jumps involve the intricate flipping of both favourable and unfavourable domains.

**Origin of orbital magnetization**

Further analysis requires understanding the origin of magnetization in MATBG. Orbital magnetization arises from Berry curvature and the topological nature of the bands, and has two contributions$^{6,13-16}$. The first is the self-rotation of the wavepacket, $M_{SR}$, and the second is the electric-field-induced transverse drift velocity of the wavepacket, which we label Chern magnetization ($M_{C}$), since it acquires its main contribution in topological Chern gaps$^{14}$. The two contributions have analogues in the quantum Hall effect, where $M_{SR}$ is similar to magnetization due to the cyclotron motion of bulk electrons and $M_{C}$ is equivalent to the magnetization induced by the equilibrium (ground-state) topological currents flowing in the edge states$^{16,40}$. It follows that $M_{C}$ is generally of the opposite sign to $M_{SR}$ (Methods).

The magnitudes and evolutions of $M_{SR}$ and $M_{C}$ in the metallic state are not universal and depend on details of the band structure. To evaluate these magnetizations, we perform single-particle band structure calculations (Methods)$^{6}$. We assume that at the charge neutrality point (CNP), the system is degenerate with $\mu_{C}$ (depending on the band structure parameters; Extended Data Fig. 2). Since no states exist in the gap, $M_{SR}$ remains constant until the next flat band with opposite valley polarization ($K^{\prime}$) is filled (Methods). Since $K$- and $K^{\prime}$-polarized flat bands have $M_{SR}$ of opposite signs, $M_{SR}$ decreases for $\nu > 1$. The corresponding $m_{SR} = dM_{SR}/d\nu$ (Fig. 4d), which is the self-rotation component of the experimentally measured total $m_{SR}$, is fairly small at low fillings and rapidly increases on approaching $\nu = 1$, reaching values of about $2 \mu_{B}$ per electron. Furthermore, $m_{SR}$ vanishes in the gap and

**Fig. 3 | Imaging local magnetization and equilibrium currents.** a, Differential magnetization $m_{SR}(x, y) = dM_{SR}(x, y)/d\nu$ reconstructed from $B_{SR}^{\pm}(x, y)$ induced by filling-factor modulation of $\nu = 0.083$ at $\nu_{c} = 0.737$. The white areas are outside the sample. b, Equivalent description of magnetization in terms of equilibrium currents, $\mathbf{J} = \nabla \times \mathbf{M}$. The magnitude of equilibrium current density $|\mathbf{J}^{\pm}(x, y)|$ reconstructed from $B_{SR}^{\pm}(x, y)$ is shown (Supplementary Video 3 shows $J^{\pm}(x, y)$ for the full range of $\nu$). c, Tomographic view of $m_{SR}(x, y, \nu)$ ($\nabla \times \mathbf{M}$). d-f, $m_{SR}(x, y, \nu)$ ($\nabla \times \mathbf{M}$) and magnetization hysteresis $\Delta m_{SR}(x, y, \nu) = (m_{SR}(x, y, \nu + 1) - m_{SR}(x, y, \nu))/2$ (f) at $\nu = 0.966$ (Supplementary Video 4 shows the tomographic view over the full range of $\nu$). The red (blue) colours indicate paramagnetic-like (diamagnetic-like) local differential magnetization. The r.m.s value of $\Delta m_{SR}$, variations in the non-hysteretic part of $\mathbf{f}$ is $0.69 \mu_{B}$ per electron, providing a measure of error of the $m_{SR}$ data.
Fig. 4 | Evolution of orbital magnetism. a, Single-particle band structure of the KA-polarized flat band in the first mini-Brillouin zone for twist angle $\theta = 1.08^\circ$, staggered potential $\delta = 17\, \text{meV}$ and tunnelling ratio $\omega = 0.95$ (Methods). b, Momentum-resolved self-rotation magnetization $m_{\text{SR}}(k)$ in the KA band. c, Calculated evolution of the momentum-integrated magnetization $M_z$ and its two components ($M_{\text{SR}}$ and $M_{\text{C}}$) versus filling factor $\nu$ in the compressible states or versus chemical potential $\mu$ in the Chern gap (cyan), which is taken to be $\Delta = M_{\text{max}} \times \mu_C = 7\, \text{meV}$ as an example. Starting from zero magnetization at the CNP, $M_{\text{SR}}$ (orange) grows positive in the KA band, remains constant in the gap and then decreases in the following K’A flat band. $M_{\text{C}}$ (purple) is very low in the compressible state and grows negative close to the band edge, followed by a linear increase with $\mu$ in the Chern gap, that is, $dM_{\text{C}}/d\mu = \mu_C/\hbar$, where $C = -1$. The total $M_z$ (dashed) changes sign in the gap region (red dot), above which an abrupt hysteretic flipping of the polarity of $M_{\text{SR}}$ and $M_{\text{C}}$ may occur due to electron recondensation from the K to K’ valley. Such flipping is not shown here for clarity and is presented in Extended Data Fig. 3. Although, theoretically, the transition through the gap should be sharp due to zero density of states, in practice, the presence of disorder substantially broadens the gap region. d, Corresponding evolution of differential magnetization $m_z$ and its $m_{\text{SR}}$ and $m_{\text{C}}$ components. Extended Data Fig. 3 provides details in the vicinity of $\nu = 1$. e–g, Tomographic slices of $\Delta m_z(y, \nu)$ (e), $m_{\text{SR}}(y, \nu)$ (f) and $m_{\text{C}}(y, \nu)$ (g) across the sample at $x = 15.2\, \mu\text{m}$. In g, the solid black line marks the first-order phase transition $\nu_f(y)$, where $m_{\text{SR}}(y, \nu)$ flips discontinuously due to electron recondensation from the K to K’ valley, coinciding with the top of the hysteretic signal $\Delta m_z(y, \nu)$ in e. The dotted black line marks $\nu_0(y)$ along which $m_z$ vanishes, demarcating the transition from the $m_{\text{SR}}$-dominated to $m_{\text{C}}$-dominated regions. A line cut along the yellow dashed line is presented in Extended Data Fig. 3b. The bottom strip schematically shows the Chern mosaic with $C = 1$ (blue) and $C = -1$ (red) patches and their valley and sublattice polarizations.
has opposite sign in the following band. The calculated $M_{\text{c}}$ behaves fairly opposite to $M_{\text{SR}}$ (Fig. 4c). It is minimum in the compressible state and attains its maximum (negative) value of $M_{\text{c}}^{\text{max}} = C\Delta e/h$ when the chemical potential $\mu$ reaches the top of the Chern insulator gap ($\Delta$ is the gap energy, $C$ is the elementary charge, $h$ is the Planck constant and $C = -1$; Fig. 4c). Note that in the metallic state, orbital magnetism depends on the evolution of the Berry curvature with chemical potential. In contrast, when $\mu$ resides in the gap, $M_{\text{c}}$ and $m_{\text{c}}$ have a universal behaviour, namely, $m_{\text{c}} = dM_{\text{c}}/d\mu = C\Delta e/h$, because the Berry-curvature flux of a full band is quantized and equal to $C$, implying finite $m_{\text{c}}$ if and only if $C$ is non-trivial. Thus, the measurement of finite $m_{\text{c}}$ constitutes a direct observation of a topological gap with the sign of $C$ given by the sign of $m_{\text{c}}$.

**Magnetcization reversal**

The K- and K'-polarized flat bands are degenerate at zero field. At low $\nu$, a small $B$ favours the valley polarization (for example, K) with positive total magnetization, that is, $M_{\text{c}} = M_{\text{SR}} + M_{\text{c}} \approx M_{\text{SR}}$. As the local chemical potential $\mu$ enters the gap, the $M_{\text{c}}$ contribution dominates and may cause $M_{\text{c}}$ to change sign (Fig. 4c, red dot). When this happens, it is energetically favourable to recondense all the states through a first-order phase transition to the opposite valley (K'), realigning $M_{\text{c}}$ with $B$. This leads to a discontinuous flip of the polarities of $M_{\text{c}}$ and $m_{\text{c}}$ and inversion of the curves$^{23,24}$ (Fig. 4c,d and Extended Data Fig. 3a).

The tomographic slice of the measured $m_{\text{c}}(y,\nu)$ (Fig. 4g) shows this flipping mechanism at play. Focusing on the $C = -1$ patches (Fig. 4g, bottom), $m_{\text{c}}$ starts off red (positive) and intensifies with $\nu$, reflecting the contribution of $m_{\text{SR}}$, reaching peak values of the order of $10 \mu_{\text{B}}$ per electron, consistent with the calculations shown in Fig. 4d. The magnetization then drops rapidly (but reversibly) through $m_{\text{c}} = 0$ (Fig. 4g, dotted line, $v_{\text{SR}}(y)$) to intense blue, reflecting the onset of negative $m_{\text{c}}$ in the Chern gap. On further filling, a discontinuous flip from blue to red occurs (Fig. 4g, solid black line, $v_{\text{SR}}(y)$), where the system recondenses all the carriers to the opposite valley. Depending on the height of the barrier between these two states, the process is hysteretic or non-hysteretic (Fig. 4e), as discussed below.

The condition for the flip to occur is $-M_{\text{SR}}^{\text{max}} > M_{\text{c}}^{\text{max}}$, where $M_{\text{SR}}^{\text{max}}$ is the value of $M_{\text{SR}}$ at the top of the band. By integrating $m_{\text{c}}(x; y,v_{\text{SR}})$ over $\nu$ up to $v_{\text{SR}}(x,y)$, we attain a map of $M_{\text{c}}^{\text{max}}(x,y)$ (Fig. 5b); $M_{\text{SR}}^{\text{max}}(x,y)$ (Fig. 5b) attains typical values of $-1 - 2 \mu_{\text{B}}$ per u.c., consistent with the calculations shown in Fig. 4c. By integrating $m_{\text{c}}(x; y,v_{\text{SR}})$ from $v_{\text{SR}}(x,y)$ over $\nu$ up to $v_{\text{SR}}(x,y)$, we attain a map of the lower bound of $M_{\text{c}}^{\text{max}}$ (Methods) (Fig. 5c). The red patches in Fig. 5b appear as dark blue/purple in Fig. 5c, with values of $M_{\text{c}}^{\text{max}}$ reaching down to $-4 \mu_{\text{B}}$ per u.c., thus explaining the mechanism of magnetization flipping. Moreover, the observation of negative local $M_{\text{c}}$ at $\nu = 1$ constitutes a direct local observation of a topological gap with a negative Chern number ($C = -1$), providing strong evidence that locally, the sample is a zero-field Chern insulator$^{34}$. From the relation $M_{\text{SR}}^{\text{max}} = C\Delta e/h$, we evaluate the lower bound of the local Chern gap; for $M_{\text{SR}}^{\text{max}} = -4 \mu_{\text{B}}$ per u.c., it is $\Delta = 7 \text{meV}$, consistent with compressibility studies$^{34}$.

**Chern mosaic**

Figure 4g shows that together with paramagnetic-like ‘red’ patches ($C = -1$) that behave as discussed above, there are numerous diamagnetic-like ‘blue’ patches ($C = 1$). These patches display an opposite sequence starting from light-blue $m_{\text{c}}$, crossing through zero to dark red, and flipping abruptly to dark blue. The presence of these patches with negative $M_{\text{SR}}$ is highly surprising since they are seemingly energetically unfavourable. Even more puzzling is that according to the described mechanism, we expect the hysteretic jumps to reflect only the abrupt flipping of negative $m_{\text{c}}$ to energetically favourable positive $m_{\text{c}}$. Remarkably, Figs. 4g and 3f show that the hysteretic jumps involve flipping of both blue and red patches.

The origin of the diamagnetic-like (blue) patches cannot be attributed to an antiferromagnetic desire to minimize the magnetic energy. The measured stray magnetic fields from neighbouring red patches are $B_{\text{SR}} \approx 1 \text{mT}$ and $B_{\text{c}} \approx 50 \text{mT}$. Thus, the energy cost of a blue patch of size $S$ ($S M_{\text{SR}} B_{\text{SR}}$) is four orders of magnitude larger than the antiferromagnetic energy gain ($S M_{\text{c}} B_{\text{c}}$), and the origin of these patches requires another explanation.

Since $M_{\text{c}} = C\Delta e/h$, the measurement of patches with opposite signs of $M_{\text{c}}$ in the vicinity of $\nu = 1$ (Fig. 4g) is a direct local observation of a mosaic of positive and negative $C$. Here $C$ is determined by the product of valley (K and K’) and sublattice (A and B) polarizations, giving four possible competing combinations. Denoting red patches (with $C = -1$) as KA polarized leaves three polarization possibilities for neighbouring blue patches: K’A, KB or K’B (ref. 23). K’B also has $C = -1$ and can be dismissed; therefore, blue patches must be polarized to either K’A or KB. Both possibilities have a magnetic energy cost ($S M_{\text{SR}} B_{\text{SR}}$) and an additional cost of either KK’ or AB domain-wall energy, proportional to the interface length. The KA–K’A mosaic is unlikely since there is no obvious external mechanism that enforces it. In contrast, the KA–KB mosaic naturally arises by sublattice polarization, which breaks the $C_{\text{sym}}$ symmetry associated with gap opening. Such a gap opening must be present in any case to form the observed zero-field Chern insulator state.

We propose two possible mechanisms to explain a space-dependent pattern of sublattice symmetry breaking, which is implied by our measurements. The first mechanism is based on Hartree–Fock calculations, which predict spontaneous sublattice symmetry breaking near $\nu = 1$, even in the absence of an external sublattice potential$^{11,12}$. If the hBN substrate potential is incommensurate with the graphene moiré lattice, it acts as a spatially varying field that couples to the local sublattice polarization. In the presence of strain and angle disorder, this field becomes effectively random, and according to the Imry–Ma argument$^{13}$, the system breaks up into domains of opposite sublattice polarizations, whose size is large in the small random field limit, thus forming a spontaneous Chern mosaic stabilized by disorder.

Alternatively, if the graphene–graphene and hBN–graphene angles create nearly commensurate moiré lattices (Methods and Extended Data Fig. 4), it has been shown$^{30,32}$ that in the absence of disorder, a position-dependent periodic sublattice polarization occurs, resulting in semimetallic or gapped Chern bands with $C$ equal to $-1$ or 0. Combining these predictions with the presence of strain$^{14}$, structural relaxation$^{15}$ and twist-angle disorder$^{16}$ in MATBG, we envision a non-homogeneous non-periodic Chern mosaic, where KA-polarized patches with positive $M_{\text{SR}}$ and $C = -1$ and KB-polarized patches with negative $M_{\text{SR}}$ and $C = 1$ are energetically stabilized by the combination of disorder and a coarse-grained substrate potential. Figure 5a shows a schematic of such a KA–KB interface at $\nu = 1$ (neglecting spin) where the bands invert across the domain wall and $C$ changes sign.

By analysing the $m_{\text{c}}(x; y,\nu)$ signal at filling factors below $v_{\text{SR}}(x, y)$ (Methods), we reconstruct the Chern mosaic structure of our sample (Fig. 5d). We find isolated $C = -1$ patches (red) of a characteristic size of a few micrometres that are embedded in $C = 1$ (blue) and $C = 0$ or semimetallic regions (green). Note that none of the phases percolate across the sample, explaining the observation of the non-quantized anomalous Hall effect in the transport (Fig. 1c–g), as proposed elsewhere$^{17}$.

**Dynamics of magnetization reversal**

In the derived picture, the A/B sublattice pattern is imprinted in the sample by local variations in the lattice structure. In the presence of small positive $B$, $\nu < 1$, the red KA patches (positive $M_{\text{SR}}$) are energetically favourable, whereas the blue KB patches (negative $M_{\text{SR}}$) are unfavourable. This raises the question of why their valley polarization does not flip to K’B to have favourable magnetization. The answer is that the $U_{\text{SR}} = -\text{mB}$ cost of the KB patches is balanced by...
avoiding the KK′ domain walls with line energy $U_{\text{nn}} \approx 0.05 \text{ meV \text{ nm}^{-1}}$ (refs. [32,42]) (Methods). Equating $U_{\text{nn}} = \pi D M_{\text{eff}} B_{y}/4 = \pi D U_{\text{nn}}$, for a patch of diameter $D$ and using $M_{\text{eff}} = 1 \mu_0$ per u.c., we attain $D_{\text{min}} = 4 l_{\text{hBN}}/M_{\text{eff}} \approx 5 \mu$m. Therefore, for a KB patch of $D < D_{\text{min}}$ it is energetically favourable to keep K-valley polarization and pay the cost of negative $M_{\text{eff}}$ rather than forming a KK′ domain wall.

Since the typical size of the Chern patches shown in Fig. 5d is $D < D_{\text{min}}$, at low $B_y$ and $\nu < 1$, the entire sample is valley polarized to K, composing a Chern mosaic of patches of different C values. On entering the Chern gap, the situation changes due to the development of large $M_z$ that is of opposite polarity to $M_{\text{eff}}$ (Fig. 4g). Now, the KA patches become energetically unfavourable (blue), and should abruptly recondense into the K′ states, as discussed above. However, individual patches cannot independently flip their valley polarization because of the high KK′ domain-wall energy, $U_{\text{nn}}$. As a result, the recondensation must occur collectively in which a large domain, composed of a number of $C = -1$ and $C = 1$ patches, reverses its valley polarization together. Note that such a large domain has smaller average $M_z$ leading to larger estimated minimum domain size $D_{\text{min}}$. Note also that KB patches, which at low $\nu$ were energetically unfavourable (negative $M_{\text{eff}}$), become energetically favourable (red) in the gap due to positive $M_{\text{eff}}$, and therefore, their flipping back to energetically unfavourable states (blue) cannot be explained without such a collective effect. Figure 4g provides direct evidence for such an abrupt collective K to K′ valley inversion where oppositely coloured patches simultaneously flip their magnetization across the black line ($\nu(x,y)$). Once a domain flips its valley polarization, a KK′ domain wall is formed along its edges, providing the means for back-gate-controllable creation of valley domain walls.

Figure 5e presents the map of the filling factor, $\nu(x,y)$, at which $m_\nu$ flips polarity on sweeping $\nu$ (Methods). It reveals that the magnetization reversal process involves the flipping of large domains with $D$ approximately equal to 5 \( \mu \)m, along with some smaller domains. There is a substantial spread in the $\nu_{\text{f1}}$ values ranging from 1.01 (right-hand side of the sample) to 1.17 (left-hand side), and some areas (yellow–grey) show no detectable flipping. On sweeping $\nu$, the $\nu_{\text{f1}}(x,y)$ map (Fig. 5f) displays a similar behaviour, although with somewhat smaller typical domain sizes. The left part of the sample shows substantially larger hysteresis in $\nu_{\text{f1}}$ in correlation with the larger $M_{\text{max}}^C$ (Fig. 5c). This can be related to the slightly higher twist angle in the left part of the sample (Extended Data Fig. 5). From $\Delta m_{\nu}(x,y,v)$, we extract the evolution of magnetization hysteresis as a function of $\nu$ (Extended Data Fig. 6), which agrees well with transport hysteresis (Fig. 1d), and explains the progressive evolution of the minor loop hysteresis (Fig. 1g).

With increasing $B_y$, $D_{\text{min}}$ should decrease and eventually drop below the Chern mosaic patch size. In this case, the magnetic energy dominates $U_{\text{nn}}$ and each patch acts independently with no collective reversals similar to the case of magnetically doped topological insulators [43]. Therefore, we anticipate that at $\nu < 1$, all the patches align with positive $M_{\text{eff}}$, forming KA and K′B patches separated by KK′ domain walls, leading to the suppression of hysteresis, consistent with transport (Fig. 1c–f).

**Outlook**

Our finding of the Chern mosaic in a sample not intentionally aligned to hBN suggests that this new type of topological disorder is ubiquitous in MATBG devices (Methods and Extended Data Fig. 4).
Moreover, identifying magnetism through transport measurements predominantly relies on the observation of hysteresis in twisted bilayer graphene. This requires a first-order transition. Since our local measurements reveal Berry-curvature magnetism and Chern gaps even in the absence of local hysteresis, orbital magnetism in MATBG could be more omnipresent than reflected by transport. Being imprinted by the substrate potential, it dictates the local flavour polarization, providing a unique mechanism for the manipulation and possible utilization of flavour-dependent electronic properties. In particular, the valley and Chern domain walls can host novel edge states. Domain walls between the regions of opposite sublattice polarizations are expected to carry two co-propagating chiral edge states, whereas domain walls where both valley and sublattice polarizations flip may carry counterpropagating valley helical modes of opposite valley polarizations. At such a domain wall, the system is likely to favour an intervalley coherent state, where the charge and valley currents are carried by a gapless collective mode. Studying the structure of these domain walls and their transport properties is an interesting direction for future investigations.

Online content
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Methods

Device fabrication. The MATBG device was fabricated using a cut-and-stack technique. All the flakes were first exfoliated on a Si/SiO2 (285 nm) substrate and later picked up using a polycarbonate/polydimethylsiloxane stamp. All the layers were picked up at a temperature of ~100 °C. The graphene was initially cut with an atomic force microscopy (AFM) tip, to avoid strain during the pickup process. The polycarbonate/polydimethylsiloxane stamp was used to pick up the top hBN and then the first graphene layer. Before picking up the second graphene layer, the stage was rotated by an angle of 1.1°–1.2°. Finally, the bottom hBN and bottom graphite gate were picked up. The finalized stack was dropped on a Si/SiO2 substrate by melting the polycarbonate at 180 °C and etched into a Hall bar using CHF3/O2.

One-dimensional contacts were formed by evaporating Cr (5 nm)/Au (50 nm). The optical and AFM images of the different fabrication steps are shown in Extended Data Fig. 7a–d. All the presented experimental data were acquired on this device that was originally studied elsewhere. The device has subsequently undergone several additional thermal cycles before the current study. A high-resolution AFM image of the device acquired after the present study is presented in Extended Data Fig. 7e. The dark spots in this image are due to damage caused to the device after the completion of the measurements presented in this paper. We do not observe any clear correlation between the topographic structure and the measured Chern mosaic.

There was no intentional alignment between the graphene bilayer and either of the hBN flakes. The white, yellow and red dashed lines in Extended Data Fig. 7a–c mark the naturally broken edges of the graphene, top hBN and bottom hBN flakes, respectively. Extended Data Fig. 7d shows that twist angle θmg between graphene and the bottom hBN layer was 6°, whereas the relative twist angle with the top hBN was 68°, which is equivalent to 8° under the symmetries of the system. These twist angles are within the range of possible angles that can create commensurability between the graphene–graphene and graphene–hBN moiré lattices (Extended Data Fig. 4).

Transport measurements. Four-point transport measurements were performed at T = 300 mK using standard lock-in techniques with a root mean square (r.m.s) bias current of I = 10 nA at 11 Hz. Figure 1a,b shows the longitudinal and transverse resistances R∥ and R⊥, respectively. The R∥ data are reproduced in Extended Data Fig. 1a along with the map of the Landau levels (Extended Data Fig. 1b). We extract the following parameters: back-gate capacitance Cbg = 450 nF cm−2, back-gate voltage VB, corresponding to v = +4 and v = −4 is 9.48 V and −9.71 V, respectively, and the carrier density for full filling of the flat band is nF = 2.69 × 1012 cm−2, which corresponds to a twist angle of θ ≈ 5.35°, where θ = 0.246 nm.

At low filling factors, the total orbital magnetization Mz = Mz∥ + Mz⊥ is dictated by Mz∥. Since K- and K'-polarized bands are equivalent, a small applied field will select the valley (K) with positive Mz∥, for the twist angles shown in Extended Data Fig. 1a–c. The transport measurements at low fields display indications of C = −1 behaviour in the vicinity of v = 1, as reported previously.

SOT fabrication and magnetic imaging. Indium SOTs were fabricated as described elsewhere, with diameters ranging from 150 to 200 nm and included an integrated shunt resistor on the tip. The SOT readout was carried out using a cryogenic SQUID series array amplifier. The magnetic imaging was performed in a ‘He system at 300 mK at which the indium SOTs can operate in magnetic fields of up to 0.6 T. At fields B∥ = 50 mT used in this study, the SOTs displayed flux noise down to 250 nΦ0 Hz−1/2, spin noise of 10μΦ0 Hz−1/2 and field noise down to 10nT Hz−1/2 at frequencies above a few hundred hertz. Since at lower frequencies, the sensitivity is limited by the 1/ν noise, the magnetic imaging was performed in the a.c. mode by applying a back-gate voltage Vbg at a frequency of 5–6 kHz. For height control, the SOT was attached to a quartz tuning fork, as described elsewhere. The tuning fork was electrically excited at a resonant frequency of ~35 kHz. The current through it was amplified using a room-temperature home-built transimpedance amplifier (designed based on another work) and measured using a lock-in amplifier. The scanning was performed at a constant height of about 160 nm above the top hBN surface. The B∥ images were acquired with a pixel size of 100 nm and acquisition time of 120 ms per pixel.

Current and magnetization reconstruction. For reconstruction of the equilibrium currents from the measured B∥ (x, y) using the invariant method described in detail elsewhere. A similar procedure was used for magnetization reconstruction. Before the inversion, a Gaussian high-pass filter was applied to the raw data to remove a weak long-range parasitic background arising from thermal gradients due to the back-gate a.c. excitation. Supplementary Video 2 shows the evolution of B∥ (x, y, t), B∥c (x, y, t) and ∆B∥+ (x, y, t) in the range of ν = 0.737–1.174. The corresponding reconstructions of m(x, y, ν) and |Fc(x, y, ν)| from B∥ (x, y, t) are shown in Supplementary Video 3. A three-dimensional tomographic representation of m(x, y, ν), m∥(x, y, ν) and ∆m∥(x, y, ν) is presented in Supplementary Video 4.

Band structure calculations. The continuum model15 is used to calculate the band structure and orbital magnetization of MATBG. Two graphene layers are coupled by the interlayer electron hopping. For each graphene layer, the lattice vectors are set as

\[ \mathbf{a}_1 = \sqrt{2} \mathbf{b}_1 = \left( \frac{1}{2}, \frac{1}{2} \right), \quad \mathbf{a}_2 = \sqrt{2} \mathbf{b}_2 = \left( -\frac{1}{2}, \frac{1}{2} \right), \]

where a1 = 0.142 nm is the intralayer bond length in graphene. The A and B sublattices are located at δx = (0, 0) and δz = (0, 1). The corresponding reciprocal lattice vectors are

\[ \mathbf{b}_1 = \frac{\sqrt{2}}{2} \left( \frac{1}{2}, \frac{1}{2} \right), \quad \mathbf{b}_2 = \frac{\sqrt{2}}{2} \left( -\frac{1}{2}, \frac{1}{2} \right), \]

and the K and K’ valleys are located at K = (±1, ±1, 0), where ζ = ±1.

The Hamiltonian is given by

\[ H = H_B + H_C + H_{bc} \]

where \( H_B \) and \( H_C \) are the top- and bottom-layer Hamiltonians, respectively, given by

\[ H_B = \sum_{q, s, \sigma} a_{s, \sigma}^\dagger (q) \mathbf{R} \sigma b_{s, \sigma} (q) \]

Here \( a_{s, \sigma}^\dagger (q) \) is the electron creation operator for an electron with spin index \( \sigma \) and momentum \( q \), and \( \mathbf{R} \sigma \) is the rotation matrix for the top/bottom layer with twist angle θ = θB = θC. The Pauli matrices, momentum \( \mathbf{q} \) is defined relative to the K valley, \( \mathbf{h} \) is the reduced Planck constant and Fermi velocity \( vF \) is taken to be 0.596 eV nm−1. The staggered potential \( \delta \), induced in bottom graphene by the hBN substrate, breaks the inversion symmetry.

The interlayer coupling is given by

\[ H_{bc} = \sum_{q, s, \sigma} \left( T_{bc} (q, \mathbf{q}) a_{s, \sigma} (q) \right) \]

Three interlayer hopping processes couple \( q \) and \( \mathbf{q} \), with momentum transfer \( \mathbf{q} - \mathbf{q} = (\mathbf{q}_A - \mathbf{q}_B - \mathbf{q}_C) \), where \( \mathbf{q}_A = \frac{2\pi}{\sqrt{3}a} (1, 1, 0) \) and \( \mathbf{q}_B = \frac{2\pi}{\sqrt{3}a} (\frac{2}{3}, \frac{1}{3}, \frac{1}{2}) \). The interlayer hopping matrix is given by

\[ T_{bc} (q, \mathbf{q}) = \frac{\delta}{w} \left( \begin{array}{cc} w & 1 & 1 \\ 1 & w & 1 \\ 1 & 1 & w \end{array} \right) \]

The staggered potential \( \delta \) is applied to the bottom graphene. As a result, the sublattice polarizations in the top and bottom layers (\( P_A \) and \( P_B \)) are not identical. They are given by the expectation value of polarization operator \( \hat{t} \) and \( \hat{n} \). In the above truncated basis, they are expressed as

\[ \hat{t}_z = \frac{1}{2} \left( \hat{c}^\dagger \hat{d} + \hat{d}^\dagger \hat{c} \right) \]

\[ \hat{n}_z = \frac{1}{2} \left( \hat{c}^\dagger \hat{c} + \hat{d}^\dagger \hat{d} \right) \]
where $\tau_i$ is the Pauli matrix in the sublattice space A/B.

Figure 4a shows the resulting structure of the KA flat band. The sublattice polarizations $P_s$ and $P_t$ are presented in Extended Data Fig. 8. Note that the sublattice polarization is only partial and is more pronounced in the bottom graphene layer to which the staggered potential is applied. The sublattice occupations weight $W_{\Lambda}$ in each layer is given by $(1 \pm P_t)/(2)$.

In this work, we have used 42 sites in the reciprocal space to construct the Hamiltonian, which gives an 84 × 84 Hamiltonian for a single flavour K. The interaction-induced Chern gap size of $\Delta = 7$ meV, comparable to the experimentally evaluated value, was manually introduced with constant $M_{\text{eff}}$ and linear $M_t$ with a slope of $dM/d\mu = C \sigma/h$, where $C = -1$.

The corresponding differential magnetization, equivalent to the signal measured in the experiment, is shown in Fig. 4d with $m_{\text{eff}} = dM/d\mu = dM/d\mu$ in the compressible regions and $m_{\text{eff}} = dM/d\mu$ and $m_{\text{eff}} = dM/d\mu$ in the gap. The following band at $r > 1$ is assumed to be of the opposite valley and thus has the same $M_{\text{eff}}$ and $m_{\text{eff}}$, evolution, but with opposite signs.

We have numerically explored the dependence of magnetization on the various parameters in the ranges of $1.00 \leq \theta \leq 1.2^\circ$, $9 \leq \delta \leq 29$ meV and $0.75 \leq \Delta \leq 1.00$.

Although the qualitative behaviour remains similar in this range of parameters, there are quantitative differences, which are particularly pronounced for variations in tunneling strength ratio $\omega$. Extended Data Fig. 2 shows the results obtained for the same $\theta = 1.08^\circ\delta = 17$ meV, but with lower $\omega = 0.8$ (ref. 39). The $m_{\text{eff}}$ and $M_t$ values are very similar, although slightly lower than those in Fig. 4c, and remain consistent with the experimental values (Fig. 5b,c). The main difference is observed in $m_{\text{eff}}$ and $M_t$, which show a more gradual drop towards negative values close to the top of the band. As a result, $m_{\text{eff}} = m_{\text{eff}} + m_t$, showing a maximum value $m_{\text{eff}}$ near $\omega = 0.9$. Extended Data Fig. 2d summarizes the values of $m_{\text{eff}}$ calculated for a range of $\omega$ and $\delta$. The simulation results for $\omega < 0.85$ are consistent with experimental values of $m_{\text{eff}} \geq 5 \mu_B$ per electron. Note that the calculations are based on the assumption of flavour degeneracy lifting from the CNP. If the degeneracy lifting occurs only above a certain filling factor $k_B^\text{F}$ (refs. 18, 30), $M_t$ will remain zero below $k_B^\text{F}$ and will raise faster above $k_B^\text{F}$, giving rise to larger $m_t$ and $m_{\text{eff}}$ on approaching the top of the band.

An important insight from the above calculations is that the Berry-curvature-induced magnetization is a sensitive probe of the local band structure. We note, however, that the calculation of magnetization was done for non-interacting electrons, and does not include the effect of interactions. Moreover, the calculations do not take into account the strain, which can be expected to further modify the magnetism. The large spatial variations in local magnetization (Supplementary Video 4) point out the complexity of the local band structure and of the substrate potential.

Evolution of $m_t$ across the KA to K′ reconfiguration phase transition. We inspect the behaviour of $m_{\text{eff}} = m_{\text{eff}} + m_t$ more closely by analysing its evolution along the dashed yellow line in Fig. 4g (Extended Data Fig. 3b). The $m_t$ value shows two sign changes as a function of $\nu$. The first occurs at $\nu = 1$ (Fig. 4g, dotted black line) that marks the transition from $m_{\text{eff}}$ to $m_{\text{eff}}$. The transition is continuous and thus $m_t$ goes smoothly through zero and is additionally smeared by our finite a.c. modulation of $\nu = 0.083$. This sign change occurs slightly below $\nu = 1$, as shown by the numerical calculations (Fig. 4d and Extended Data Fig. 3a), and its position depends on the band structure parameters.

The second sign change occurs at $\nu = 1$ (Fig. 4g, solid black line) due to a completely different mechanism of carrier recombination from valley K to K′ through a first-order transition. As shown in Fig. 4c, the total magnetization $M_t$ decreases with increasing the chemical potential in the gap region due to the negative contribution of $M_t$. The $M_t$ is dominated and thus $m_{\text{eff}}$ dominates magnetization. This transition is continuous and thus $m_t$ goes smoothly through zero and is additionally smeared by our finite a.c. modulation of $\nu = 0.083$. This sign change occurs slightly below $\nu = 1$, as shown by the numerical calculations (Fig. 4d and Extended Data Fig. 3a), and its position depends on the band structure parameters.

Note that in the absence of recondensation transition, for $\nu = 1/2$, the carriers first fill the KA band followed by filling of the K′ band. In the presence of a transition, however, the sequence is reversed: the carriers first fill the KA band followed by filling of the K′ band. The simulation results for $\nu = 1/2$ show the development of a discontinuous flattening of magnetization. Figure 4d shows the expected evolution of magnetization in the absence of this flattening transition, whereas Extended Data Fig. 3a presents the calculated $m_t$ value with flattening. Since this transition is discontinuous and hysteretic, it is not smeared by the finite $\nu$ and thus remains discontinuous in the experimental data (Fig. 4g, black line) and seen in the jump between the neighbouring experimental data points (Extended Data Fig. 3b).
\[ \theta_{\text{BIN}}^\pm = \arccos \left( \frac{r - 1}{\sqrt{(r - 1)^2 + s^2}} \right) \]

where \( r = 1.017 \) is the ratio between the hBN and graphene lattice constants, \( s = (r - 1)^2 + (\frac{q}{p} + \frac{p}{q})/2 \), and \( q/p \) and \( p/q \) are a triplet of co-prime integers. The resulting super-moiré cell is \( \sqrt{r} \) times larger than the \( m_z \) moiré cell; thus, we only consider low values of \( n = 1 \) and 2. Solutions to the above equations are shown in Extended Data Fig. 4 for values of \( \theta_{\text{BIN}} \) close to the magic angle. Numerous additional solutions exist for larger \( n \). Note that due to the symmetry of the system, each \( \theta_{\text{BIN}}^+ \) and \( \theta_{\text{BIN}}^- \) solution has six-fold rotational symmetry. Furthermore, it is shown that the \( \theta_{\text{BIN}}^+ \) and \( \theta_{\text{BIN}}^- \) can vary by some amount from the commensurate angle (depending on the value of \( n \)), implying that there is a finite width around each solution that will lead to a Chern mosaic (Extended Data Fig. 4, error bars). Additionally, it was pointed out elsewhere that the commensurate angles probably create a lower-energy state, implying that a sample might relax into these states. Finally, taking into account the possible commensurability with either top or bottom hBN, the presence of strain, charge disorder, \( \theta_{\text{BIN}}^\pm \) approaches and presumably \( \theta_{\text{BIN}} \), the formation of moiré commensurability might be ubiquitous. In particular, \( \theta_{\text{BIN}} \) values of 8° and 6° of the top and bottom hBN, respectively, (Extended Data Fig. 7, optical images) are within the range of possible angles for the formation of commensurate moiré lattices (Extended Data Fig. 4).

Note that the above analysis ignores all types of disorder and specifically twist-angle disorder that is known to be substantial in MATBG samples. A more realistic picture includes twist-angle disorder in both hBN and graphene, which has higher sensitivity at low fields but precludes the imaging of twist-angle disorder that is known to be substantial in MATBG samples. A more realistic picture includes twist-angle disorder in both hBN and graphene, which has higher sensitivity at low fields but precludes the imaging of twist-angle disorder that is known to be substantial in MATBG samples. A more realistic picture includes twist-angle disorder in both hBN and graphene, which has higher sensitivity at low fields but precludes the imaging of twist-angle disorder that is known to be substantial in MATBG samples. A more realistic picture includes twist-angle disorder in both hBN and graphene, which has higher sensitivity at low fields but precludes the imaging of twist-angle disorder that is known to be substantial in MATBG samples. A more realistic picture includes twist-angle disorder in both hBN and graphene, which has higher sensitivity at low fields but precludes the imaging of twist-angle disorder that is known to be substantial in MATBG samples. A more realistic picture includes twist-angle disorder in both hBN and graphene, which has higher sensitivity at low fields but precludes the imaging of twist-angle disorder that is known to be substantial in MATBG samples. A more realistic picture includes twist-angle disorder in both hBN and graphene, which has higher sensitivity at low fields but precludes the imaging of twist-angle disorder that is known to be substantial in MATBG samples. A more realistic picture includes twist-angle disorder in both hBN and graphene, which has higher sensitivity at low fields but precludes the imaging of twist-angle disorder that is known to be substantial in MATBG samples.

Experimental derivation of Chern mosaic. Differential magnetization \( m_x(x,y,v) \) is a continuous function of \( v \), except at the abrupt flipping of magnetization due to the recondensation of carriers from K to K' valley. We are interested in determining the filling factor \( v(x,y) \), where \( m_x(x,y) \) continuously traverses zero from \( m_x(x,y) \) dominated to \( m_x(x,y) \)-dominated magnetization as \( v(x,y) \) varies from the Chern gap, and the filling factors \( v(x,y) \) and \( v(x,y) \) where \( m_x(x,y) \) discontinuously flips its sign. The finite width of our a.c. modulation of the filling factor (\( \Delta v \)) peak to peak) causes some averaging of \( v(x,y) \) samples. In another study, a map of twist angle \( \theta \) was measured over the entire range of filling factors. Note that averaging has no essential effect on \( v(x,y) \) angles probably create a lower-energy state, implying that a sample might relax

\[ \Delta \theta = \arccos \left( \frac{r - 1}{\sqrt{(r - 1)^2 + s^2}} \right) \]

from the transport data). The derived local twist angle along this line, that is, \( \Delta \theta \), causes some averaging of \( v(x,y) \) samples. In another study, a map of twist angle \( \theta \) was measured over the entire range of filling factors. Note that averaging has no essential effect on \( v(x,y) \) angles probably create a lower-energy state, implying that a sample might relax

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Competing interests
The authors declare no competing interests.

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Correspondence and requests for materials should be addressed to Eli Zeldov.

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Extended Data Fig. 1 | Landau fan diagram. a, Longitudinal resistance $R_{xx}$ at $T=300$ mK reproduced from Fig. 1a. b, Landau level fits. Integer fillings $\nu$ of the moiré supercell are indicated by vertical grey lines and the Landau levels and Chern insulating states are indicated by diagonal lines along with their corresponding Chern numbers. Dotted lines are extrapolations.
Extended Data Fig. 2 | Orbital magnetization calculations. a, The evolution of the integrated magnetization $M_z$ (dashed) and its two components $M_{SR}$ and $M_C$ vs. the filling factor $\nu$ in the compressible region and vs. $\mu$ in the gap region (cyan) in the KA flat band for twist angle $\theta = 1.08^\circ$, staggered potential $\delta = 17$ meV, and tunnelling ratio $w = 0.8$. b, The corresponding evolution of the differential magnetization $m_z$ and its $m_{SR}$ and $m_C$ components. c, The self-rotation magnetization $m_{SR}(k)$ in the first mini-Brillouin zone in the KA flat band. d, The calculated peak value of the differential magnetization $m_{peak}^z$ near the top of the band as a function of $\delta$ and $w$ for $\theta = 1.08^\circ$. 
Extended Data Fig. 3 | Evolution of $m_z$ through recondensation phase transition. 

**a.** Numerically calculated $m_z$ in the vicinity of $\nu = 1$, similar to Fig. 4d, but including the first-order recondensation transition marked by the red dot, upon which the $m_z$ flips its sign discontinuously. The value of the Chern gap is taken to be 3.5 meV and $m_z = dM_z/dn$ is calculated taking into account the finite $v^n = 0.083$ modulation. 

**b.** The measured $m_z$ along the vertical yellow dashed line in Fig. 4g. The black dots are the experimental data points and the solid curve is guide to the eye. The first sign change of $m_z$ at $\nu \approx 0.9$ marks the continuous crossover from $m_{SR}$ to $m_c$ dominated magnetization. The second sign change at $\nu \approx 1$ is discontinuous (and hysteretic) and reflects the first-order $K$ to $K'$ recondensation transition.
### Extended Data Fig. 4 | Moiré commensurability conditions

Calculated pairs of graphene–graphene $\theta_{GG}$ and graphene–hBN $\theta_{GBN}$ twist angles that result in commensurability of the two moiré lattices following Ref. 20. Each solution is labelled by its $(n, p, q)$ integer triplet. Shown are only solutions for $n = 1, 2$ as the super moiré cell scales like $n^2$, but numerous additional solutions exist for larger values of $n$ with proportionally larger super moiré cells. The blue error bars show deviations from the exact commensurate angles where a Chern mosaic structure is still expected to form.
Extended Data Fig. 5 | Local twist angle. \(a\), \(B_z^{2D}(x)\) measured along the dotted line in Fig. 2b vs. the global filling factor \(\nu\) in the vicinity of \(\nu = 4\) (top) and \(\nu = -4\) (bottom). The black lines mark \(\nu^+\) and \(\nu^-\) corresponding the local electron and hole dispersive band edges, \(n^+\) and \(n^-\). \(b\), The derived local twist angle \(\theta(x) = \sqrt{\frac{1}{2} n^+ - n^-}\).
Extended Data Fig. 6 | Comparison of magnetization and transport hysteresis. a, Percentage of sample area that has locally hysteretic $m_z$ represented as a histogram as a function of filling factor $\nu$. b, The hysteresis in transport, $\Delta R_{yx}(\nu) = R_{yx}(\nu_\uparrow) - R_{yx}(\nu_\downarrow)$, at $B_s = 47$ mT from data presented in Fig. 1d.
Extended Data Fig. 7 | Optical and AFM images of the device. a–d, Optical images of graphene (a), top hBN (b), bottom hBN (c) and the final stack (d). The white (graphene), yellow (top hBN), and red (bottom hBN) dashed lines show that the top (bottom) hBN layer is at an angle of 68° (6°) with respect to the graphene. e, Zoomed-in AFM image of the device. The two large dark spots are due to damage caused to the device after the measurements presented in this paper.
Extended Data Fig. 8 | Sublattice polarization. a–b. Single particle band structure of the K valley with colour indicating the degree of sublattice polarization $P$ in the bottom (a) and top (b) graphene layers. Red (blue) colour shows polarization of A (B) sublattice with $P = 0.2$ corresponding to 60% occupation weight on A and 40% on B sublattices. The sign of $P$ determines the sign of C. Staggered potential $\delta = 17$ meV is applied to the bottom graphene, $\theta = 1.08^\circ$ and $w = 0.95$. 