Quantum Hall Effect of Massive Weyl Fermions in \( n \)-type Tellurene Films

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The advent of graphene has evoked the re-examination of band topology of Dirac/Weyl nodal materials which can host low-energy realistic quasiparticles. Under strong magnetic fields, the topological properties of two-dimensional Dirac/Weyl materials can be directly manifested through quantum Hall states. Here we report the first observation of massive Weyl fermions through quantum Hall effect in n-type Weyl semiconductor tellurene (two-dimensional form of tellurium). The n-type doping profile forms a wide quantum well in tellurene, where two correlated layers of electrons create a pair of symmetric-antisymmetric energy states in addition to spin and valley degeneracy, leading to an approximate SU(8) isospin symmetry. The chirality-induced Weyl nodes residing near the edge of the conduction band give rise to radial spin texture, and topologically non-trivial $\pi$ Berry phase was detected in quantum Hall sequences. Our work presents strong evidence of massive Weyl fermions and expands the spectrum of Weyl matters into semiconductor regime.
The emergence of graphene\(^1\) has triggered vast interest to study the properties of realistic particles in low energy spectrum of topological materials. For example, under high magnetic fields, unconventional sequence of graphene quantum Hall states at filling factor \(\nu = 4(n + \frac{1}{2})\) (here \(n = 0, 1, 2 \ldots\)) suggests \(\pi\) Berry phase of Dirac fermions\(^2,3\), which is a direct consequence of the topological nature of Dirac nodes. Since then, many other classes of topological materials with Dirac/Weyl nodal features in band structure have been predicted and identified with great potentials for spintronics, optoelectronics, and quantum computing applications. However, these Dirac/Weyl nodal points created by the crossing of either two bands or two branches of spin polarized bands are generically limited to semimetal category without exploitable bandgaps. Here we introduce a new type of topological matters - Weyl semiconductor tellurene (two-dimensional form of tellurium), where the Weyl nodes are located at the edge of conduction band, allowing us to investigate the properties of Weyl fermions in the context of semiconductors.

Trigonal tellurium (Te) is a narrow bandgap semiconductor with a unique one-dimensional helical atomic structure. The crystal consists of three-fold screw-symmetric atomic chains interconnected by van der Waals-like forces\(^4\), as shown in Fig. 1a. The screw symmetry distinguishes two irreducible enantiomers with opposite chirality in real space, which falls into either P3\(_2\)12 or P3\(_1\)21 space group, depending on the handedness. As a consequence of lacking inversion symmetry and heavy atoms, tellurium shows strong spin-orbit coupling (SOC) effect. One prominent feature of the band structure occurs at the bottom of the conduction band at H point (Fig. 1b and 1c), where the spin degeneracy is lifted to form band crossing due to strong SOC, which is often referred to as a “Rashba-like” band\(^9,10\). This is because the band dispersion is similar to typical parabolic band with the
presence of Rashba SOC term, while the hedgehog-like radial spin texture is distinctive from Rashba bands\textsuperscript{10} (as we shall further discuss later). This crossing point between two spin polarized bands at the conduction band edge (see Fig. 1c) can be classified as a Weyl point\textsuperscript{10–13}, which gives rise to massive Weyl fermions with quantum topological properties similar to recently proposed Kramers-Weyl nodes in chiral crystals with strong SOC\textsuperscript{14}. Note that in Te since H point is not a time reversal invariant momenta (TRIM) point, the degeneracy is not protected by Kramers theorem (which is the prerequisite of Kramers-Weyl nodes), but instead guaranteed by its three-fold rotational symmetry. Nonetheless, despite the different mechanism leading to Weyl nodes, they both share similarity in band topology and spin textures.

To date, most of the magneto-transport measurements were performed on bulk Te samples\textsuperscript{15–17}, where the resolution of Shubnikov-de Haas (SdH) oscillation features were limited by the sample quality as well as weak Landau quantization. Recently, a hydrothermal growth method is proposed, yielding atomically flat two-dimensional (2D) Te films with dangling-bond-free surfaces\textsuperscript{5,6,8} (Fig. 1e inset), with thickness ranging from a couple of atomic layers to tens of nanometers. These 2D Te films are named tellurene solely for simplicity\textsuperscript{5}, as an analogy to other 2D elemental materials including graphene, stanene and phosphorene. Much better developed SdH oscillations as well as quantum Hall states were observed in these p-type tellurene films\textsuperscript{18}, adding a new member into the scarce family of high-mobility 2D materials hosting quantum Hall effect (QHE) including graphene\textsuperscript{2,3,19}, black phosphorous\textsuperscript{20–22}, InSe\textsuperscript{23}, and some transition metal dichalcogenides (TMDs)\textsuperscript{24–27}. However, due to the unintentional p-type dopants in Te, the chemical potential is usually fixed near the valence band edge. Therefore so far most of the
experiments are limited to holes and little about the properties of the conduction band has been investigated via transport measurements\textsuperscript{6,18,28}. Here we managed to dope tellurene films into n-type using atomic layer deposited (ALD) dielectric doping technique without degrading electron mobility\textsuperscript{29,30}, which grants us the access to the conduction band and explore much more exotic physics of massive Weyl fermions. The schematic of n-typed doped device is demonstrated in Fig. 1d. The as-synthesized tellurene films were dispersed on to a degenerately doped silicon substrate with a 90 nm SiO\textsubscript{2} insulating layer on top, followed by patterning and deposition of Ti/Au metal contacts. Here we chose low work function metal contacts Ti to reduce the electron Schottky barrier height and accommodate n-type transport. A layer of alumina was subsequently deposited onto tellurene film by ALD grown at 200 °C, converting the tellurene film underneath from p-type to n-type, as confirmed by $I_d-V_g$ transfer curves of a tellurene field-effect transistor (FET) at both room temperature and cryogenic temperature (see Fig. 1e). Similar ALD doping method has also been reported on other material systems like black phosphorus\textsuperscript{31–33} and silicon\textsuperscript{34}. The doping mechanism is attributed to the threshold voltage shift caused by positive fix charges in low-temperature ALD-grown films\textsuperscript{31} or the interface electric dipole field\textsuperscript{34}.

For magneto-transport measurements, six-terminal Hall bar devices were fabricated as shown in Fig. 2a (see Methods for more fabrication and measurement details). We investigated over 20 devices with typical film thickness ranging from 10 to 20 nm, and all of them exhibit similar and reproducible behaviors in general. Here all the data presented is from one high-quality device, unless otherwise specified. The global back gate allows us to tune the 2D electron density from 2 to 10$\times$10\textsuperscript{12} cm$^{-2}$. Representative longitudinal ($R_{xx}$) and transverse ($R_{xy}$) magnetoresistance curves measured at $V_g=10$V are plotted in Fig. 2b,
with Hall density of $2.5 \times 10^{12}$ cm$^{-2}$ and Hall mobility of 6,000 cm$^2$/Vs. The onset of SdH oscillations is around 2 T, leading to an estimated quantum mobility of 5,000 cm$^2$/Vs, which is close to Hall mobility within a reasonable margin. At B field around 24 T and 32 T, $R_{xy}$ is fully quantized into integer fraction of $h/e^2$ (corresponding to filling factor of 3 and 4), and $R_{xx}$ also drops to 0 -- a hallmark of QHE$^{35}$. As shown in Fig. 2c, by fixing B field at 42 T and sweeping back gate, all filling factors from 2-8 are resolved (although not all fully quantized due to Landau level broadening), suggesting all the degeneracies, including spin and valley, have been lifted and the system reaches quantum Hall ferromagnetism regime. By mapping out $R_{xx}$ through the $V_g$-B parameter space, we can construct the Landau fan diagram in a color map as in Fig. 2d. The data section from 0 to 12 T and from 12 T to 45T were acquired from two different magnet systems at 30 mK and 300 mK, respectively, therefore the data shows slight discontinuity at B=12 T. Since the conduction band edge is located at two inequivalent H (H’) points in the first Brillouin zone accounting for two-fold valley degeneracy, it is intuitively conceivable that each Landau level consists of four degenerate energy states (2 for valley degeneracy and 2 for spin degeneracy) like graphene, with cyclotron energy gap $E_C = \hbar \omega_c = \frac{\hbar e B}{m^*}$ at filling factor of 4, 8, 12… Here $\hbar$ is the reduced Planck constant, $e$ is the electron charge, and $m^*$ is the effective mass. Following this argument, we should expect the energy gap to increase monotonically in the sequence of 4n as we approach lower filling factors, since the cyclotron energy $\frac{\hbar e B}{m^*}$ increases linearly with the magnetic field. Yet we notice from Fig. 2b that the gap at $\nu = 12$ is larger than that at $\nu = 8$, suggesting the origin of these two gaps are different and the four-level single particle picture cannot explain these unconventional sequence, hence another degree of freedom needs to be taken into account.
To understand this anomaly, we first focus on SdH oscillation features at relatively low B field regime (see Fig. 3a). When B field is below approximately 7 T, only one set of oscillations is observed, and the oscillation frequency $B_F$ can be used to extract carrier density: $n_{SdH} = eB_F/h = 1.02 \times 10^{12} cm^{-2}$, where degree of degeneracy is not taken into account. Meanwhile we can also calculate the total carrier density from the slope of $R_{xy}$ to be $n_{Hall} = 8.75 \times 10^{12} cm^{-2}$, which is approximately 8 times larger than $n_{SdH}$, implying that there is eight-fold degeneracy. At B field over 7 T, two sets of oscillations are resolved as marked by red and blue arrows alternatively in Fig. 3a. There are several plausible explanations of two sets of oscillations: (a) a second sub-band; (b) another electron pocket at other point(s) of Brillouin zone; (3) two coupled layers of charge at the top and bottom of a wide quantum well\(^3_{36-39}\). However, both multiple sub-bands and pockets assumptions involve two Fermi surfaces and introduce two sets of independent oscillations with different frequencies since $B_F$ is proportional to the area of the Fermi surface, and the ratio of two Fermi surface areas can be modulated by tuning the carrier density. This will be reflected in two independent oscillation frequencies in the fast Fourier transform. However, in our experiments only one frequency (and its higher order of harmonics) can be found in frequency domain (see Supplement Information 1), and the ratio of $n_{Hall}/n_{SdH}$ remains to be 8 throughout the entire gate range that we can probe (Fig. 3b). This solidifies the coupled wide quantum well assumption to be the only reasonable explanation. Because of the doping profile of ALD dielectrics, two layers of carriers are induced at the top and the bottom of the Te film when appropriate gate bias is applied (Fig. 3c top panel). Should the central potential barrier be low enough, electron tunneling is permitted between two layers, and their wave functions will interfere and be reconstructed into symmetric and anti-
symmetric states (bottom panels of Fig. 3c and 3d). In this case, the Landau levels in two layers will become resonant despite a small energy gap $\Delta E_{SAS}$, which is related to the width of the quantum well (film thickness) and applied gate voltage. When $\Delta E_{SAS}$ is sufficiently small compared to Landau level gaps (as shown in the inset of Fig. 3a), we can treat the near degenerate symmetric-antisymmetric energy levels as another type of isospin which, along with real spin and valley isospin, yields an approximate SU(8) symmetry group$^{40}$. To further verify the wide quantum well assumption, we fabricated a double-gate device by adding a metal gate on top of ALD alumina, as shown in Fig. 3c. $R_{xx}$ is mapped out as a function of both top gate voltage and back gate voltage at a fixed B field of 7 T in Fig. 3e. The tilted strips in the lower left corner with negative slope suggest that two layers of electrons are coupled and the Fermi level can be tuned by both of the gates. However when both gates are strongly positive biased (in the upper right corner of Fig. 3e), the electrons are strictly localized near two surfaces and each gate can only modulate one set of oscillations. In this case the tilted lines will evolve into straight lines, which is a signature of wide quantum well and has been observed in other quantum well systems$^{41,42}$. Ideally, in the upper right regime of Fig. 3e, we should observe a “chessboard” feature with two sets of oscillations each controlled by one gate independently$^{42}$. However the horizontal patterns are almost indiscernible, probably due to the low mobility at top interface between tellurene and ALD alumina.

We now focus on assigning the quantum Hall isospin ferromagnetic states to each filling factor by fixing the perpendicular magnetic field $B_\perp$ and increasing the total magnetic field $B_{total}$. Three competing mechanisms can break SU(8) isospin symmetry: (1) the real spin can be resolved by Zeeman splitting with an energy gap proportional to the total
magnetic field: \( \Delta E_Z = g^* \mu_B B_{\text{total}} \), where \( g^* \) is effective g-factor and \( \mu_B \) is the Bohr magneton; (2) \( \Delta E_{SAS} \) is in general irrelevant to perpendicular magnetic field \( B_\perp \), but the in-plane magnetic field can destroy this energy gap and collapse the symmetric and antisymmetric energy levels\(^{36} \); (3) the valley isospin is usually polarized by breaking inversion symmetry or electron interactions\(^{43,44} \). Fig. 4b shows \( R_{xx} \) versus \( V_{bg} \) at fixed \( B_\perp = 25 \, \text{T} \) and \( B_{\text{total}} \) is increased from 25 T to 45 T. The gaps in \( \nu=5 \) and \( \nu=7 \) start to collapse with increasing \( B_{\text{total}} \), which can be attributed to the in-plane magnetic field destroying \( \Delta E_{SAS} \)\(^{36} \) (another plausible explanation of band crossing due to large Zeeman splitting like in WSe\(_2\) case\(^{45} \) can also be ruled out by absence of coincident effect, see Supporting Information 2). Therefore we can assign the odd filling factors to the symmetric-antisymmetric energy gap \( \Delta E_{SAS} \). On the contrary, the gaps in filling factor \( \nu=6 \) and \( \nu=8 \) are expanded, however this can be related to either the increasing of the Zeeman energy gap, or the collapsing of \( \Delta E_{SAS} \) in neighboring odd filling factors. Hence, at this point we cannot explicitly determine the sequence of spin and valley splitting, and further investigation is needed to completely understand the mechanism that drives the system to fully polarized ferromagnetism regime.

Finally we present the evidence of massive Weyl fermions in conduction band edge of Te. Weak anti-localization (WAL) effect is observed in near-zero magnetic field regime (see Supporting Information 3), manifesting strong SOC effect in Te system -- a direct consequence of lacking inversion symmetry and heavy atoms of Te. The strong SOC gives rise to spin-polarized band crossing at H and H’ point (as shown in Fig. 5a), which is protected by three-fold rotational symmetry. This point at H (H’) can be classified as a Weyl node that can be viewed as a Berry curvature monopole in the momentum space. We shall address that unlike normal band-inversion-induced Weyl points which usually lead to
linear band dispersion of Weyl semimetals that can hardly be tuned by gate, the chirality-induced Weyl nodes in Te are located at the edge of the conduction band, and therefore Te belongs to a new type of topological material-- Weyl semiconductor, which grants us more freedom to probe the topological properties of Weyl nodes and to design novel electronic and quantum devices by taking advantage of the versatility of the semiconductors. Evidence of Weyl fermions such as negative magnetoresistance and kinetic magnetoelectric effect in bulk Te has been reported. Here we present a much more convincing evidence-- the observation of non-trivial π Berry phase in magneto-transport. The amplitude of SdH oscillations, when neglecting all degree of degeneracy, is described as: \( \Delta R_{xx} \propto R(B,T) \times \cos (2\pi(B_F/B + 1/2) + \varphi) \), where \( B_F \) is the oscillation frequency and \( \varphi \) is the Berry phase. To extract the Berry phase, we record the values of \( 1/B \) at each \( \Delta R_{xx} \) minima and plot them against the Landau level index, as shown in Fig. 5b. The extrapolated the linear fitting curves intercept with y-axis at 1/2 (see Fig. 5c and 5d), which is the signature of non-trivial π Berry phase, as observed in many other Dirac or Weyl topological materials. To assertively rule out any potential offset from imperfection of the device, we measured over 10 devices and all of them unanimously show near 1/2 intercept (see Supporting Information 4). The origin of the Berry phase is substantially rooted in the hedgehog-like spin texture near the chirality-induced Weyl nodes. The spin texture are either pointing at or away from the Weyl nodes as illustrated by red arrows in Fig. 5a. Under magnetic field, when an electron completes a cyclotron motion in the real space, its momentum changes by \( 2\pi \), corresponding to a closed loop trajectory around the Fermi surface in momentum space. Along this path, the spin of the electron also rotates by \( 2\pi \), which picks up a π Berry phase, since electrons are spin-1/2
particles. The origin of Berry phase in Te resembles that near the Dirac points of graphene, except that in graphene the Berry phase is induced by radial pseudospin (valley isospin) texture rather than the real spin as in Te.

Band-inversion-induced nodal structures are usually accompanied with linear band dispersion that gives rise to the massless realistic particles like Dirac/Weyl fermions. However this is not necessarily always the case. The Dirac/Weyl nodes only guarantee the topological properties, but the band dispersion can still be arbitrary, depending the energy scale of interest. Here we measure SdH oscillation amplitudes as a function of temperature (Fig. 6a), which can be described by the Lifshitz–Kosevich equation: \[ \Delta R_{xx} \propto \frac{2\pi^2 m^* k_B T}{\sinh (2\pi^2 m^* k_B T / e B)} \]. The effective mass is then extracted to be about 0.10 m\(_0\) (where m\(_0\) is the bare electron mass) throughout the entire gate range, which is consistent with previously reported value of bulk Te\(^4\,52\). This suggests that although the Weyl nodes reside only several meV above the conduction band edge, the topological properties (Berry phase) and the band mass carry over to a much broader energy window of at least 50 meV as the gate can access to. Therefore n-type tellurene film is an ideal playground to study the behavior of massive Weyl fermions with tunable chemical potentials.

In conclusion, using ALD dielectric doping on high-quality tellurene films, for the first time, we observed well-developed quantum Hall effect from 2D electrons in tellurene and accessed the electronic structure of its conduction band. A wide quantum well model with two correlated electron layers is proposed to explain the anomaly in SdH oscillations and quantum Hall sequences. Eight-fold degeneracies, including real spin, valley isospin and symmetric-antisymmetric isospin are all lifted under a 45 T external magnetic field, leading to fully polarized quantum Hall ferromagnetic states. Topologically non-trivial \( \pi \) Berry
phase is ambiguously detected as a direct evidence of massive Weyl fermions in the vicinity of chirality-induced Weyl nodes with predicted hedgehog-like spin textures. Our work offers a material platform to explore the topological properties of massive realistic quasiparticles in chiral semiconductors with strong SOC.
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Figure 1 | Crystal structure of Te and n-type ALD doped tellurene device. a, Crystal structure of Te with helical atomic chains. b, First Brillouin zone of Te. The conduction band minima is located at the corner points H and H’. c, Zoomed-in dispersion of lowest spin-split bands along K-H-A line. A Weyl node is formed at H point, as highlighted by the red circle. d, Schematic view of an n-type Te device with global back gate and ALD top doping layer. e, I_d-V_g transfer curves of a typical ALD doped n-type device measured at room temperature and cryostat temperature. Inset: An optical image of an as-synthesized tellurene film obtained from hydrothermal growth method. The helical chains (z-axis) are aligned along the longer edge of the film, as indicated by the red arrow.
Figure 2 | Quantum Hall effect in tellurene two-dimensional electron gas. a, An optical image of a six-terminal Hall bar device. b, Longitudinal ($R_{xx}$, in blue) and transverse ($R_{xy}$, in red) resistance as a function of magnetic field. c, $R_{xx}$ and $R_{xy}$ a function of back gate voltage. The gate oxide is 90 nm SiO$_2$. d, Color mapping of $R_{xx}$ by sweeping both back gate voltage and magnetic field. The data from 0-12T and 12-45T are measured from the same device in a superconducting magnetic system at 30 mK and a hybrid magnetic system at 300 mK, respectively.
Figure 3 | Eight-fold degeneracy of Landau levels in tellurene coupled wide quantum well. a, SdH oscillations evolving from eight-fold degeneracy to two sets of resonant oscillation features (blue and red arrows). Inset: illustration of degenerate energy levels with symmetric-antisymmetric energy splitting. b, Comparison of Hall density and SdH density under different gate biases. c, Schematics of double-gated devices and carrier distribution of separate layers (top) and correlated layers (bottom). d, Potential profile and wave function distribution of separate layers (top) and correlated layers (bottom). e, Color mapping of normalized $R_{xx}/R_{xx}(B=0)$ by sweeping both top gate and back gate. The data is acquired at a magnetic field of 7 T.
Figure 4 | Quantum Hall ferromagnetic states under a titled B field. a, R$_{xx}$ versus gate voltage measured at a fixed perpendicular magnetic field 25 T and increasing total magnetic fields (with 100 $\Omega$ offset). Here the dashed lines correspond to filling factors from 4 to 8. b, Sequence of competing mechanism lifting degeneracies within a Landau level.
Figure 5] Weyl nodes and Berry phase near Te conduction band edge. a, Weyl node at H point induced by strong SOC in chiral crystal. The red arrows around the Fermi surface represent hedgehog-like spin texture. b, Landau fan diagram under different gate bias. The scattered symbols are read off from the value $1/B$ of each minima in $R_{xx}$ and plotted against the Landau level index. Straight lines are linear fitting under each gate bias. c, Zoom in of black box in b with linear fittings extrapolated to the y-axis. d, Intercept of linear fittings versus gate voltage. The error bars come from the linear fitting. The reference line at $y=0.5$ corresponds to $\pi$ Berry phase.
Figure 6| Temperature-dependent SdH oscillations and effective mass of Weyl fermions. a, SdH oscillation amplitudes (subtracting a smooth background) under various temperature from 0.5 K to 18 K. Inset: $\Delta R_{xx}$ versus temperature fitted by Lifshitz–Kosevich equation. b, The effective mass is extracted to be about $0.10 m_0$, and it is independent on gate biases and magnetic fields.
Online Content Methods, along with any additional Extended Data display items are available in the online version of the paper; references unique to these sections appear only in the online paper.

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Supporting Information for:

Quantum Hall Effect of Massive Weyl Fermions in n-type Tellurene Films

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Methods

Synthesis of 2D tellurium crystals

In a typical procedure\(^1\), analytical grade Na\(_2\)TeO\(_3\) (0.00045 mol) and certain amount of poly(-vinyl pyrrolidone) was resolved into deionized water (33 ml) at room temperature under magnetic stirring to form a homogeneous solution. The resulting solution was transferred into a Teflon-lined stainless-steel autoclave, which was then filled with aqueous ammonia solution (25\%, w/w\%) and hydrazine hydrate (80 \%, w/w\%). The autoclave was sealed and maintained at the reaction temperature for a designed time. Then the autoclave was cooled to room temperature naturally. The silver-gray, solid products were precipitated by centrifuge at 5000 rpm for 5 minutes and rinsed 3 times with distilled water (to remove any ions remaining in the final product).

Langmuir-Blodgett (LB) transfer of 2D Te films

The hydrophilic 2D Te nanoflake monolayers can be transferred to silicon substrates by the Langmuir-Blodgett (LB) technique\(^2\). The washed nanoflakes were suspended in a mixing solvent of N, N-dimethylformamide (DMF) and CHCl\(_3\) (e.g., in the ratio of 1.3:1). Then, the mixture was dropped into the deionized water. After the evaporation of the solvent, 2D Te flakes will be floating on the surface of water. Then we can scoop 2D Te films onto the substrates.

Device fabrication

The as-grown 2D Te films were dispersed onto heavily doped Si substrates with 90 nm SiO\(_2\) insulating layer, followed by DI water rinse and standard solvent cleaning process
(acetone, methanol and isopropanol). Hall-bar devices were patterned with electron beam lithography (EBL) and 30/90 nm Ti/Au metal contacts were deposited with electron beam evaporator under the pressure below $2 \times 10^{-6}$ Torr. To eliminate the geometry non-ideality, the device was then trimmed into standard Hall bar shape with better symmetry using BCl$_3$/Ar dry etching. 20 nm of Al$_2$O$_3$ was deposited onto Te films at 200 °C with atomic layer deposition for n-type doping. For double-gated device, another Ni/Au top metal gate was patterned and deposited with EBL and electron beam evaporator subsequently to cover the entire channel region.

**Magneto-transport measurements**

Part of the low magnetic field transport measurements were performed a Triton 300 (Oxford Instruments) dilution fridge system with 12 Tesla superconducting coils. The high magnetic field data were acquired in a 31 Tesla resistive magnet system (Cell 9) and a 45 Tesla hybrid magnet system (Cell 15) in National High Magnetic Field Lab (NHMFL) in Tallahassee, FL. The electrical data was recorded using standard low-frequency AC measurement technique using SR830 lock-in amplifiers. The transfer curves of two-terminal FETs with a Cascade probe station and Keysight B1500A semiconductor analyzer at room temperature and current mode of lock-in amplifier at 30 mK.
Supporting Information 1: Fast Fourier transform of SdH oscillations

There are several common scenarios where two sets of SdH oscillations can be observed. For instance, (a) two sub-bands of 2D quantum wells are involved when the Fermi level fills up to the second sub-band; (b) another pocket of carriers in Brillouin zone exists with energy level close to the conduction band edge, and (c) symmetry-antisymmetry states in a wide quantum well. Since the SdH oscillation frequency is proportional to the area of the Fermi surface, we should be able to distinguish two sets of independent oscillations by performing fast Fourier transform (FFT) for (a) and (b) assumption. Here we performed FFT on $R_{xx} vs 1/B$ under various gate bias, and only one frequency (and its higher orders of harmonics) can be observed, suggesting this feature in SdH oscillations is caused by two sets of resonant energy levels.
**Figure S1**| Fast Fourier transform of SdH oscillations under different gate bias. Only single frequency and its higher harmonics can be detected.

**Supporting Information 2: Shubnikov-de Haas oscillations in tilted magnetic field**

As a common practice to extract effective g-factor, SdH oscillations are often measured in tilted magnetic field to disentangle the cyclotron energy $E_C = \frac{\hbar e B_\perp}{m^*}$ and Zeeman energy $E_Z = g^* \mu_B B_{total}$. By rotating the sample to certain angle, one should expect to see the coincidence effect -- the SdH maxima and minima will exchange when the criteria $E_C = E_Z$ is met. Here we measured our sample in tilted magnetic field, however coincidence effect is not observed up to 78.2° (see Fig. S2), from which we can estimate the upper bound of effective g-factor to be 2.68. The real $g^*$ can be even much smaller given the fact that the no clear trace of developing coincidence effect is observed up till 78.2°. This suggests that the Zeeman energy in Te is much smaller compared to cyclotron energy and we can safely rule out the possibility that the gap closing in Fig. 4a is caused by the spin-split level crossing as in the case of WSe$_2$.3.
**Figure S2**| SdH oscillations in tilted magnetic field. No evidence of coincidence effect is observed, suggesting a small effective g-factor.

**Supporting Information 3: Weak anti-localization in n-type Te films**

Weak anti-localization (WAL) is the phenomena that at low B field regime, the resistance of the material shows a dip centered at zero. In normal material there should be a peak in resistance at zero field, since the magnetic field and destroy the localization effect and reduce the resistivity of the material, which is referred to as weak localization effect. However in systems with strong spin-orbit coupling, the trajectory of the electrons travelling around the disorder clockwise and counterclockwise will contribute the opposite
sign and the wave functions become constructive which induce the reverse effect. The WAL of p-type Te films has been reported in the previous work\textsuperscript{4,5}. Here we show some preliminary data of WAL in n-type Te films in Fig. S3 to verify strong SOC effect in Te system. More comprehensive analysis of n-type Te WAL effects is beyond the scope of this work and will be reported independently.

**Figure S3**| Weak anti-localization of n-type Te films.

**Supporting Information 4: More data on extracting Berry phase**

The way we adopted to extract Berry phase relies on accurate reading of $R_{xx}$ minima at each filling factor. Here we only use the SdH oscillation periods where the degeneracy is not lifted. However, even the splitting is not yet resolvable at low field, there is a chance
that the energy splitting within the Landau level can deviate $R_{xx}$ minima and cause misinterpretation in extracting Berry phase. To eliminate this uncertainty, we extracted the intercepts of Landau fan diagram for another 8 devices as shown in Fig. S4, and all of them are close to $1/2$. This is a solid proof that this phenomena is reproducible and our method of extracting Berry phase is reliable.
Figure S4 | Additional data of extracting Berry phase from other eight devices.
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