Giant and anisotropic many-body spin–orbit tunability in a strongly correlated kagome magnet

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Owing to the unusual geometry of kagome lattices—lattices made of corner-sharing triangles—their electrons are useful for studying the physics of frustrated, correlated and topological quantum electronic states1–9. In the presence of strong spin–orbit coupling, the magnetic and electronic structures of kagome lattices are further entangled, which can lead to hitherto unknown spin–orbit phenomena. Here we use a combination of vector-magnetic-field capability and scanning tunnelling microscopy to elucidate the spin–orbit nature of the kagome ferromagnet Fe₃Sn₂ and explore the associated exotic correlated phenomena. We discover that a many-body electronic state from the kagome lattice couples strongly to the vector field with three-dimensional anisotropy, exhibiting a magnetization-driven giant nematic (two-fold-symmetric) energy shift. Probing the fermionic quasi-particle interference reveals consistent spontaneous nematicity—a clear indication of electron correlation—and vector magnetization is capable of altering this state, thus controlling the many-body electronic symmetry. These spin-driven giant electronic responses go well beyond Zeeman physics and point to the realization of an underlying correlated magnetic topological phase. The tunability of this kagome magnet reveals a strong interplay between an externally applied field, electronic excitations and nematicity, providing new ways of controlling spin–orbit properties and exploring emergent phenomena in topological or quantum materials10–12.

Understanding and manipulating correlated quantum materials are prerequisites for exploring their potential for applications10–12, and quantum materials that exhibit a giant response in the presence of an external field are particularly promising10–12. Kagome antiferromagnets are central in the search for exotic quantum states because both the spin and the charge are frustrated geometrically, enabling the formation of spin-liquid phases and topological electronic structures1–9. However, the realization of such states in real materials has been limited. Kagome ferromagnets are also of great interest because their unusual physics can be probed in transport, such as in transition-metal stannides13–18. Transport measurements in this family have demonstrated large anomalous Hall effects that can arise from non-trivial electronic topology with non-vanishing Berry curvatures13–18. The Berry phase of the antiferromagnet Mn₃Sn is associated with a non-collinear spin texture and the existence of topological fermions in its band structure15–17. For the soft ferromagnet Fe₃Sn₂, it is speculated that the Berry phase is associated with a massive Dirac band near the corner of the Brillouin zone, which hosts a two-dimensional gap18. Accordingly, this family serves as a fertile platform for exploring the interplay between magnetism and quantum electronic structure in kagome lattices. Here we study the atomically resolved electronic structure of Fe₃Sn₂ at 4.2 K (Curie temperature, T_{Curie} = 670 K) by using a combination of time-reversal-breaking vector-magnetic-field capability and high-resolution scanning tunnelling microscopy (STM) and spectroscopy (STS). Although many previous works focused on the unusual transport properties of Fe₃Sn₂, we observe an unexpected giant anisotropic vector-field response of the electronic states of the kagome lattice, which opens up the opportunity to demonstrate controlled quantum-level manipulation of an exotic topological phase. The methodology described here offers a new way of discovering magnetic topological phases in a strongly correlated setting, which can be used for the discovery of other correlated topological materials.

Fe₃Sn₂ has a layered rhombohedral crystal structure with space group R3m and hexagonal lattice constants a = 5.3 Å and c = 19.8 Å. It consists of a honeycomb Sn layer sandwiched between kagome FeSn bilayers (Fig. 1a). Owing to the weak bonding of these bilayers, the sample tends to cleave with either a FeSn- or a Sn-terminated surface. These two surfaces are identified experimentally via comparisons of their respective step edge heights in the crystal structure (Fig. 1b). Mapping the differential conductance of these two surfaces also reveals differences in the electronic structure (Fig. 1c). A detailed inspection of Fig. 1d confirms the honeycomb lattice structure of the Sn surface, whereas the FeSn surface exhibits smaller corrugation, hindering direct atomic identification. By analysing the line-cuts across the step edge of the Sn surface (Fig. 1e), we determine the position of the Sn atom that corresponds to the centre of the Sn honeycomb unit and those of the Fe atoms on the FeSn surface (Fig. 1d).

Having identified the two surfaces, we study the quasiparticle excitations under the perturbation of an external magnetic field. At zero field, the low-energy differential conductance spectra of both surfaces show high-intensity states around the Fermi energy, with the spectrum of the FeSn surface exhibiting an additional state at ~8 mV. Increasing the c-axis field causes a pronounced shift of the side peak towards negative energies, whereas there is no discernible shift of the states near the Fermi energy. From the observed magnetic response and its surface dependence, it is likely that the side peak arises from the magnetic Fe orbital in the kagome lattice. We find that the magnetic field response of the lattice extends beyond the Zeeman effect in several aspects. The shift of the side peak saturates around 1 T, with a total energy shift of 12 meV (Fig. 2a). We observe an identical shift when the 1 T field is reversed. More importantly, the saturation behaviour agrees well with the magnetization curve (Fig. 2c), denoting a magnetization-driven shift. If this energy shift was attributed to the Zeeman effect, it would amount to an anomalously large g factor, which has not been reported in the literature (Fig. 2b).

To explore the magnetization response of Fe₃Sn₂ in three dimensions, we rotate the external field in the a–b plane. We find that this state of the FeSn surface also saturates before 1 T when the field is applied.
in plane and that the saturated shift at 1 T evolves with the azimuth angle, θ (Fig. 2d). In contrast to the six-fold crystal symmetry of Fe₃Sn₂, this evolution has a two-fold symmetry, which can be described by the two-fold-symmetric function as \(3.2 - 3.2 \cos(2\theta)\) meV (Fig. 2e). Notably, there is no shift when the field is applied along the a axis \((\theta = 0)\). As the net magnetization is known to lie in plane at low temperatures in zero field, such nodal energy shifts lead to an effective magnetic field of around 210. From the initial shift rate below 1 T, we derive an effective g factor of around 210. The error bars are based on the energy resolution. The magnetization correlates strongly with the energy shift. The inset shows that the field is applied perpendicular to the sample surface. Dependence of the differential spectra of the FeSn and Sn surfaces on a magnetic field parallel to the c axis. The data can be fitted by a two-fold-symmetric function as \(3.2 - 3.2 \cos(2\theta)\).

The inset illustrates the field’s azimuth angle with respect to the a axis of the kagome lattice. The error bars are based on the energy resolution.

**Fig. 2 | Vector-magnetization-induced giant and nematic energy shift.** a, Dependence of the differential spectra of the FeSn and Sn surfaces on a magnetic field parallel to the c axis. The spectra are offset for clarity and the red arrows mark the peak positions. b, Energy shift of the electronic state of the FeSn surface as a function of the c-axis field. From the initial shift rate below 1 T, we derive an effective g factor of around 210. The error bars are based on the energy resolution. c, Bulk c-axis magnetization curve, expressed in units of the Bohr magneton, \(\mu_B\), per Fe atom. The magnetization correlates strongly with the energy shift. The inset shows that the field is applied perpendicular to the sample surface. d, Dependence of the differential spectra of the FeSn and Sn surfaces on the angle of an in-plane field \((B = 1\, \text{T})\). The blue circles mark the peak positions. e, Energy shift as a function of azimuth angle, \(\theta\). The data can be fitted by a two-fold-symmetric function as \(3.2 - 3.2 \cos(2\theta)\). The inset illustrates the field’s azimuth angle with respect to the a axis of the kagome lattice. The error bars are based on the energy resolution.
behaviour can be understood by considering the spontaneous magnetization to be along the $a$ axis, which already saturates the energy shift. The anisotropic evolution that we observe in our STM data also agrees qualitatively with the bulk transport anisotropy in response to a vector field (Extended Data Fig. 3), consistently demonstrating the existence of electronic nematicity in FeSn$_2$, which has been previously observed in correlated materials$^{21-24}$.

To further study the symmetry of the electronic state realized in this material, we map the differential conductance of the FeSn surface over a large area under various vector-field conditions. By taking their Fourier transforms, we obtain the quasiparticle interference (QPI) data for the electron scattering involving the band structure. The QPI data obtained at the energy of this side-peak state are shown in Fig. 3. The zero-field QPI data in Fig. 3a exhibit ring-like signals at the larger wave vectors ($q_2$) and a two-fold pattern around the zone centre ($q_1$). This spontaneously broken symmetry state also aligns with the sample’s $a$ axis, in agreement with the aforementioned nematicity, which is consistent with our transport data. Remarkably, we find that the $c$-axis magnetization removes the electronic nematicity and restores the rotational symmetry (Fig. 3b). Although the $a$-axis magnetization retains the same nematic pattern at $q_1$ (Fig. 3c), this pattern is systematically rotated by the rotation of in-plane magnetization. When the external field is withdrawn, the nematicity recovers to that shown in Fig. 3a, regardless of magnetization history. This suggests that there exists an intrinsic nematic order pinning the spontaneous magnetization along the $a$ axis. By contrast, the QPI around $q_2$ remains approximately isotropic regardless of the field, indicating that the shifting of the electronic state is probably associated with the states near $q_1$.

Such an association is further supported by the field-dependent QPI dispersion plotted in Fig. 3d, e. The QPI dispersions in Fig. 3d shows a clear hole-like band ($q_2$) with no discernible field dependence, corresponding to the ring-like signal seen in the QPI images in Fig. 3a–c. On the other hand, at small $q$, the dispersion reveals a maximum in the QPI intensity at approximately the energy of the side peak both at zero field ($E = -10 \pm 2$ mV) and at $B = 1$ T ($E = -20 \pm 2$ mV). This QPI intensity extends to larger $q$ with increasing energy, suggesting an electron-like band with a band bottom at this low-$q$ maximum-energy peak ($q_1$). Moreover, when the energy window of the dispersion is extended to even lower energies (Fig. 3e), a second hole-like band appears at $B = 1$ T, with two (one upper and one lower) branches forming a non-monotonically dispersing feature with an hourglass-like shape. The non-monotonic shape of the observed signal indicates that the scattering is sensitive to the details of the underlying dispersion in the band structure. It is consistent with a massive Dirac-like dispersion, under the assumption that the scattering is intra-band in nature and corresponds to the Dirac feature expected around this energy in the photoemission measurement$^{18}$. These two dispersive branches move farther apart from the centre of the hourglass-like feature as the
magnetic field is tuned to zero. Recent photoemission-based identification of a band structure resembling massive Dirac fermions in this material suggests a gap size of about 30 mV, consistent with our data at $B = 0$ T in Fig. 3e, despite what appears to be a shift in the chemical potential, which is possibly due to surface termination effects or doping differences in the samples. It is well known that photoemission measurements lack the ability to probe the field dependence of this gap (mass of the Dirac band) or determine whether the Dirac fermions originate from the FeSn kagome lattice or the Sn honeycomb lattice, which is critical for understanding and correctly modelling the state realized in this material.

A summary of our experimental findings is shown in Fig. 4. Our results demonstrate a vector–magnetization–based energy shift of the quantum electronic states with an intriguing correspondence to symmetry breaking (Fig. 4a, b). These states form an electron band crossing the Fermi level (Fig. 4c). Without an external field, spontaneous magnetization is along the $a$ axis; the band bottom (identified as the side peak in the tunnelling conductance in Fig. 2a) exhibits a QPI with a two-fold symmetry. Although the symmetry of the QPI rotates with the magnetization, indicating strong spin–orbit coupling (SOC) that intertwines the orbital space with the magnetic space, it is unexpected that the energy of the band bottom modulates substantially with the angle of rotation. The observed energy shift induced by in-plane magnetization also has a two-fold symmetry with its nodal line along the $a$ axis (Fig. 4a), which is indicative of an intrinsic nematic order that pins the spontaneous magnetization direction and leads to this energy difference. Rotating the magnetization to the $c$ axis causes the largest energy shift (Fig. 4a).

These giant electronic responses driven by the magnetization direction go well beyond Zeeman physics and point to a spin–orbit-entangled, correlated magnetic topological phase, which we discuss below.

In fact, previous STM work on other systems has shown that owing to the presence of SOC, the electronic structure of magnetic thin films with domain walls and skyrmions can depend on the spin orientation. Because electronic structures on the kagome lattice have linear band-crossing Dirac points at the Brillouin zone corners, it is natural for us to consider a picture of Dirac fermions in the presence of SOC in the study of kagome-lattice (quantum) anomalous Hall materials. The observed energy shift should thus result from the interplay of the Dirac gap (Fig. 4c) with magnetism. The large ferromagnetic moment splits the Dirac crossing into two branches that are well separated in energy, with spins polarized parallel and anti-parallel to the direction of magnetization. In ref. 18, the magnetization is assumed to be along the $c$ axis, and a Kane–Mele-type SOC that preserves the spin component $S_z$ is considered to produce a Dirac gap. However, the spontaneous magnetization in Fe$_3$Sn$_2$ lies in plane at low temperatures, so that the $S_z$ SOC cannot generate a gap at the Dirac crossing in the kagome lattice. This contradicts our observation of the largest mass gap (smallest energy shift) for $a$-axis magnetization, sketched in Fig. 4a, c. Thus, the physics governing the interplay between SOC and magnetism here lies beyond the Kane–Mele scenario. One possibility is that all SOC interactions respecting the full crystal symmetry need to be constructed with both $S_z$- and in-plane $S_{xy}$-conserving terms, and our results indicate that the latter should have a larger effect. Alternatively, the Dirac gap may have an additional source that interferes with that due to a dominant Kane–Mele SOC. Because Fe$_3$Sn$_2$ displays a large anomalous Hall effect with skyrmion excitations, it is possible that the latter has a contribution from the spin Berry phase associated with chiral spin textures. The spin chirality produces a gauge flux and, according to the theory, opens a Dirac gap independent of the magnetization direction. As the magnetization is rotated to the $c$ axis, the orbital flux induced by the Kane–Mele SOC can be out of phase and compete with the gauge flux, leading to the reduction of the Dirac gap (Extended Data Fig. 11), consistent with our interpretation of the data.

Furthermore, our experiment reveals an intriguing nematic order in this kagome magnet. In addition to the magnetization-controlled charge nematicity due to SOC effects, there exists an intrinsic nematic order originating from the charge channel, as evidenced by the anisotropic energy shift and transport response to the vector magnetization, as well as the pinning of the spontaneous magnetization direction irrespective of vector magnetization history. Interestingly, the well known
intra-unit-cell (q = 0) charge-ordered state on the kagome lattice, which is driven by inter-site Coulomb interactions, is a nematic state, as demonstrated theoretically.\textsuperscript{29,30}

In summary, our experiment uncovers a correspondence between vector-field-based energy shift and broken symmetry in Fe\textsubscript{2}Sn\textsubscript{2}, which demonstrates unusually large and anisotropic magnetic tunability in a spin–orbit kagome magnet and points to an underlying correlated magnetic topological ground state. The novelty of this work is the spin–orbit tunability and the gigantic response of the kagome material, which are not implied by, or can be derived from, known transport or photoemission effects. The gigantic spin–orbit response that we discovered in this strongly correlated material is unexpected and not implied by results reported in refs\textsuperscript{13,14,18}. Our findings collectively show the rich and unconventional physics of kagome magnets, which encompasses entangled magnetic, charge and orbital degrees of freedom, as well as symmetry breaking and topological properties of electronic states involving low-energy fermions. A complete understanding of this physics would require a comprehensive quantum many-body theory that describes electrons in the kagome lattice in the presence of strong spin–orbit coupling. Our space–momentum exploration of electronic excitations by controlled vector-field manipulation is a powerful tool for probing the physics of topological matter beyond weakly interacting $Z\textsubscript{2}$ topological insulators.\textsuperscript{27,28}

**Online content**

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**Author contributions**

J.-Y.X. and S.S.Z. conducted the STM and STS experiments in consultation with M.Z.H.; H.L., W.W., C.X. and S.J. synthesized and characterized the sequence of samples; K.J., G.C., B.Z., B.L., T.-R.C., H.L., Z.-Y.L. and Z.W. carried out theoretical analysis in consultation with J.Y. and M.Z.H.; I.B., T.A.C., H.Z., S.-Y.X. and G.B. contributed to sample characterization and instrument calibration; J.-Y.X., S.S.Z. and M.Z.H. performed the data analysis and figure development and wrote the paper with contributions from all authors; M.Z.H. supervised the project. All authors discussed the results, interpretation and conclusion.

**Competing interests**

The authors declare no competing interests.

**Additional information**

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METHODS

STM measurements. Single crystals of Fe₃Sn₂ with sizes of up to 1 mm × 1 mm × 0.3 mm were used in this study. Samples were cleaved mechanically in situ at 77 K in ultrahigh-vacuum conditions and then immediately inserted into the STM head, which was already at He–4 base temperature (4.2 K). Before applying the magnetic field, the tip was withdrawn to about 10 μm from the sample. The vector magnetic field was applied with the zero-field cooling technique, after which we carefully approached the tip to the sample to find the same atomic–scale area to perform tunnelling spectroscopy. Tunnelling conductance spectra were obtained with an Ir–Pt tip using standard lock-in amplifier techniques with a root-mean-square oscillation voltage of 0.5 mV and a lock-in frequency of 973 Hz. The conductance maps were taken with the tunnelling junction set up \( V = -50 \) mV and \( I = 200 \) pA, and the tunnelling spectra were taken with the junction set up \( V = -100 \) mV and \( I = 0.8 \) nA.

Sample preparation. Single crystals of Fe₃Sn₂ were synthesized by the Sn–flux method with a molar ratio of Fe:Sn = 1:19. Fe and Sn grains were placed in a clean and dry alumina crucible. Then the alumina crucible was sealed in a tantalum tube with Ar at a pressure of about 0.7 bar, which restrained the volatilization of Sn during the heating process. Finally, the tantalum tube was sealed in a quartz tube with an Ar environment at 2 mbar. The quartz tube was placed in a furnace and heated to 1,150 °C from room temperature, kept at 1,150 °C for 48 h, cooled to 910 °C in 6 h and then cooled to 800 °C at a rate of 1.5°C h⁻¹. The quartz tube was moved quickly into the centrifuge to remove the excess Sn flux from Fe₃Sn₂ single crystals. The shape of a single crystal is hexagonal with shiny surfaces.

Transport and magnetic measurement. The longitudinal resistivity \( \rho_{xx} \) and interlayer (along c axis) resistivity \( \rho_{cc} \) were measured using a standard four-probe method on a Quantum Design physical properties measurement system. The magnetic moment was measured using a Quantum Design magnetic property measurement system.

Extended Data Fig. 1 shows the longitudinal resistivity, \( \rho_{xx} \), measured for temperatures from 380 K to 5 K. The residual resistivity ratio, defined as \( \rho_{xx}(300 \text{ K})/\rho_{xx}(5 \text{ K}) \), was calculated to be 58.4. This large number attests to the high quality of our single crystal studied here.

Extended Data Fig. 2 compares the in-plane and out-of-plane magnetization curves measured at 5 K. Although both curves show a saturated value close to 2/\( \mu_B \) per Fe atom, the in-plane magnetization saturates much sooner, suggesting that the easy magnetization axis lies within the \( a-b \) plane. This is in agreement with our vector-field STM measurements, which showed that the spontaneous magnetization is along the \( a \) axis.

Extended Data Fig. 3 shows the evolution of the interlayer resistivity as a function of the azimuth angle of an in-plane field. Consistent with the vector-field STM study, it also reveals a two-fold-symmetric evolution with the azimuthal angle of the field, demonstrating intrinsic electronic nematicity.

STS measurement. As can been seen from Extended Data Fig. 4, strong interferences arise from native surface defects as either dark or bright spots in the topographic image. The bright spots are probably Sn adatoms and the dark spots are probably either Sn or Fe vacancies (see Extended Data Fig. 5 for details regarding these assignments), both of which can be generated by incomplete cleavage at low temperature. The vectors magnetic field, so the Dzyaloshinskii–Moriya interaction is not directly relevant.

Comparison with the non-atomic cleavage surface of Mn₃Sn single crystals. Although the weaker interlayer bonding between the honeycomb Sn layer and the FeSn bilayer allows the preferred cleaving planes to exist in Fe₃Sn₂, the Mn₃Sn bonds along the \( c \) axis and within the \( a-b \) plane are almost equivalent (similar to the situation within the FeSn bilayer), leaving no natural atomic cleaving plane. This is also suggested by the three-dimensional crystal shape of Mn₃Sn and two-dimensional crystal shape of Fe₃Sn₂, as shown in Extended Data Fig. 8a–8b. We have also successfully cleaved Mn₃Sn single crystals and imaged flat surfaces. However, they all exhibit non-atomic structures, as shown in Extended Data Fig. 8c–8d, for example. These images are in sharp contrast to those measured on Fe₃Sn₂, indicating that Fe₃Sn₂ is ideal for performing state-of-the-art STM studies in the transition-metal stannide family.

Theoretical discussion of the magnetization-direction-dependent Dirac gap. The model of massive Dirac fermions in a ferromagnetic kagome metal describes the Fe₃Sn₂ band structure obtained by ARPES rather well. However, this description depends on the direction of the net magnetization. For simplicity, we can first ignore the interlayer coupling, which does not change the conclusion. The effective Hamiltonian can be written as

\[ H = H_0 + H_{SO} + H_m \]

where \( H_0 \) represents the tight-binding model of the kagome lattice with nearest-neighbour hopping, \( t \):

\[ H_0 = \sum_{ij} t_{ij} c_{i}^{\dagger} c_{j} \]

Here \( c_{i}^{\dagger} (c_{i}) \) is the electron annihilation (creation) operator in the spinor notation.

\[ H_{SO} = \frac{1}{2} \sum_{ij} \lambda_{ij} \vec{s}_{i} \cdot \vec{s}_{j} \]

where \( \lambda \) is the SO coupling constant. The effective SO coupling is

\[ H_{SO} = \sum_{ij} \lambda_{ij} \vec{s}_{i} \cdot \vec{s}_{j} \]

where \( \vec{s}_{i} \) is the spin operator for the Dirac fermions.
where $\lambda$ is the SOI amplitude, $s_i$ is the spin Pauli matrix and $v_{ij} = 2(\mathbf{d}_i \times \mathbf{d}_j) \cdot z/\sqrt{3}$, with $\mathbf{d}_i$ and $\mathbf{d}_j$ denoting the unit vectors along the two bonds that the electron traverses from site $i$ to site $j$ on the kagome lattice, as shown in Extended Data Fig. 9. We note that the Kane–Mele SOI only contains spin $s_i$ without spin-flipping, thus, the $z$ component of the spin is conserved.

$H_m$ is the double-exchange coupling, $J_{H}$, between the ferromagnetic moment $\mathbf{m}$, and the conduction electrons $3$:

$$H_m = -J_H \sum_i \mathbf{s}_i \cdot \mathbf{m}_i.$$

Clearly, when $J_{H}$ or $\mathbf{m}$, is zero, $H$ generates the massive Dirac fermions as discussed in ref. $28$. Because Fe$_3$Sn$_2$ is ferromagnetic with a large ferromagnetic moment, the energy scale of $J_{H}$ $\mathbf{m}$ can be large relative to the SOI $\lambda$. The magnetization $\mathbf{m}$ splits the spin-degenerate bands into the spin minority and majority bands, with their separation controlled by $J_{H}$ and $\mathbf{m}$.

By diagonalizing the Hamiltonian $H$, we find that the Dirac mass gap depends on the direction of $\mathbf{m}$, as shown in Extended Data Fig. 10. When the magnetization is $\mathbf{m}_i = m z$, that is, along the $z$ direction ($M_z$), $H_{SO}$ causes scattering of the electrons within the spin minority and majority sectors individually. This opens the Dirac mass gap, as shown in Extended Data Fig. 10a by the red lines. On the other hand, if $\mathbf{m}_i$ is in the $x$-$y$ plane ($M_{xy}$), $H_{SO}$ scatters the electrons between the well separated spin minority and majority sectors. Thus, $H_{SO}$ cannot open the Dirac mass gap, as shown in Extended Data Fig. 10a by the blue lines. Experimentally, the magnetization direction of the ferromagnetic order in Fe$_3$Sn$_2$ is along the $a$ direction in the $x$-$y$ plane. In this sense, $H$ cannot explain the gap opening observed in ARPES. The same conclusion is reached in the case when interlayer coupling is included, where the band structure is shown in Extended Data Fig. 10b.

Because the Kane–Mele SOI cannot open the Dirac mass gap when the polarization ferromagnetic moment lies in the $a$-$b$ plane, the description of the electronic structure in Fe$_3$Sn$_2$ must go beyond the Kane–Mele approach. There are at least two possibilities, as discussed in the main text. The first one is to consider other types of or more complete spin–orbit coupling terms based on the full crystal symmetry. Such a microscopic description is difficult at the present time. Another possibility, which we briefly describe here, is that the ferromagnetic order in Fe$_3$Sn$_2$ is non-coplanar with a non-zero spin chirality $3$. In the large $J_{H}$ limit, the conduction electron spin is forced to align with the spin direction at each site. Then, the hopping terms acquire a Peierls phase resulting from rotating the electron spin to the local spin direction. This gives rise to a spin Berry phase $\Omega_{ab}$.

As shown in Extended Data Fig. 11a, the spin chirality $\chi_{ab} = \mathbf{S}_i \cdot (\mathbf{S}_j \times \mathbf{S}_k)$ corresponds to the spin Berry phase $\Omega_{ab}$, which is equal to half the solid angle formed by the spins $\mathbf{S}_i$, $\mathbf{S}_j$ and $\mathbf{S}_k$. An electron hopping around this spin configuration acquires a gauge flux $\Omega$ and opens the Dirac gap on the kagome lattice $3$. Interestingly, this mass gap is independent of the direction of the net magnetization. By including both the Kane–Mele SOI and the spin Berry phase, the effective model in the local spin basis in the spin minority sector can be written as

$$H_{\text{eff}} = \sum_v t_e^{\alpha\beta} f_i^\dagger f_j^\alpha + i\lambda_\alpha(\theta)f_i^\dagger f_j^\alpha$$

for $t \ll \lambda$, where $f_i$ is the fermion operator in the rotated basis and the effective SOI strength $\lambda_\alpha(\theta)$ depends on the net magnetization direction measured by the angle $\theta$ away from the plane. When the magnetization lies in the plane, $\lambda_\alpha(\theta) = 0$, while when it is along the $z$ direction, $\lambda_\alpha(\theta) \propto \lambda$. As a result, the spin Berry phase from a non-coplanar ferromagnetic order opens the Dirac gap while $\lambda$ does not contribute to the gap opening in the absence of an external field. When an external magnetic field rotates the saturated ferromagnetic order to the $z$ direction, the SOI becomes important and interferes with the spin Berry phase $\Omega$. As shown in Extended Data Fig. 11c, when the magnetization lies in the $x$-$y$ plane ($M_{xy}$), $\Omega$ dominates and opens the Dirac gap. Furthermore, the Dirac gap can become smaller ($\lambda \Omega > 0$) when the magnetization is along the $z$ direction ($M_z$), consistent with our interpretation of the experimental data. It is also possible that the Dirac gap for in-plane magnetization arises from purely orbital mechanisms (such as the effects of orbital hybridization discussed in ref. $22$ and staggered orbital order $34,15$), which competes with the Kane–Mele SOC when magnetization is rotated to lie along the $c$ axis $26$. However, given the existence of large anomalous Hall and related Berry-curvature effects established by transport measurements in this material, we adopt a Berry-curvature interpretation, which can then consistently describe both transport and spectroscopic data.

**Data availability.** The data that support the findings of this study are available from the corresponding author on reasonable request.

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Extended Data Fig. 1 | Longitudinal resistivity measurement. The longitudinal resistivity $\rho_{xx}$ was measured from 380 K to 5 K at zero magnetic field. The current was applied along the $a$-$b$ plane.
Extended Data Fig. 2 | Magnetization at 5 K in different magnetic field directions. The sharp saturation of the magnetization curve for the in-plane field suggests that the easy magnetization axis lies in the $a-b$ plane.
Extended Data Fig. 3 | Interlayer resistivity evolution as a function of in-plane field azimuth angle. The black line is the fitting curve. The inset shows a schematic of the sample (black hexagon) with the concentric electrical contacts used in such measurements (inner blue circle, voltage contacts; outer blue ring, current contacts).
Extended Data Fig. 4 | Spectroscopic maps with vector-field conditions.

a, Topography of a large FeSn surface. b, Differential conductance map taken in the area shown in a at the energy of the side peak. c, Differential conductance maps with various vector-field conditions taken in the area shown in a at the energy of the side peak. The black arrow indicates the a axis and the red arrow represents the field vector. Their corresponding FFT images (known as QPI signals) are shown in Fig. 3.
Extended Data Fig. 5 | Topographic image of a FeSn surface, showing two kinds of dark-spot defect. The centres of the defects are located at Fe and Sn sites, based on the atomic assignment from Fig. 1.
Extended Data Fig. 6 | Nonmagnetic scattering channels for helical spin–momentum texture with different magnetization directions.

**a**, The upper panel shows that spins cant towards the c axis to generate a net c-axis magnetization. The lower panel shows that backscattering is allowed in all directions.

**b, c**, To generate a net in-plane magnetization, the momentum of the band shifts perpendicular to the in-plane direction and the spins reorganize their directions (similar to the Edelstein effect31), as shown in the upper panels (the faded red arrows represent helical spin texture without magnetization for reference). The spins are still locked perpendicular to the original momentum centre31 (black circle). The lower panels show that backscattering is allowed only in the magnetization direction and forbidden in the perpendicular direction owing to spin reversal.
Extended Data Fig. 7 | Lorentz transmission electron microscopy images of the magnetic domain structures taken at different temperatures after turning off the external 0.7-T magnetic field applied along the \( c \) axis. The measurement is taken in the \( a\text{--}b \) plane in these images and the different shades indicate the different orientation of the spin relative to the \( a\text{--}b \) plane. The skyrmions and stripe domains gradually disappear when the temperature drops to 130 K.
Extended Data Fig. 8 | STM results on Mn₃Sn. a, b, Single crystals of Mn₃Sn and Fe₃Sn₂, respectively. c, d, Typical STM images of the cleavage surface of Mn₃Sn showing no clear atomic lattice structure.
Extended Data Fig. 9 | Kagome lattice model. a, Kagome lattice with Kane–Mele SOI. Here, \( v_{AB} = 2(d_1 \times d_2) \cdot z/\sqrt{3} = -1 \), where \( d_1 \) is the unit vector from A to C and \( d_2 \) is the unit vector from C to B. b, Flux \( \Omega \) enclosed by the fundamental plaquette formed by sites A, B and C on the kagome lattice.
Extended Data Fig. 10 | Band dispersion from the kagome model. a, Red lines are the band structure for \( M \) along the \( z \) direction (\( M_c \)) with a Dirac mass generated by \( \lambda \). Blue lines are the band structure for \( M \) in the \( x-y \) plane (\( M_{ab} \)), showing a vanishing Dirac mass gap even in the presence of \( \lambda \).

b, Similar to a but including bilayer splitting due to interlayer coupling. Only the spin minority bands are shown.
Extended Data Fig. 11 | Effect of spin Berry phase. 

a, The spin Berry phase $\Omega_{ijk}$ is equal to half the solid angle formed by spins $S_i$, $S_j$ and $S_k$. 
b, An electron hopping in the chiral spin background is equivalent to hopping in the presence of a gauge flux $\Omega$. 
c, Band structures in the presence of both $\lambda$ and chiral flux $\Omega$, when the magnetic field and the magnetization are along the x–y plane ($M_{ab}$) and the z direction ($M_c$).