Manipulation of Dirac band curvature and momentum-dependent $g$ factor in a kagome magnet

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The Zeeman effect describes the energy change of an atomic quantum state in a magnetic field. The magnitude and direction of this change depend on the dimensionless Landé $g$ factor. In quantum solids, the response of the Bloch electron states to the magnetic field also exhibits the Zeeman effect, with an effective $g$ factor that was theoretically predicted to depend on the momentum$^{1,2}$, and which may be particularly strong in topological and magnetic systems. However, the momentum dependence of the $g$ factor is difficult to extract and it has not been directly measured. Here we report the experimental discovery of a momentum-dependent $g$ factor in the kagome magnet YMn$_6$Sn$_6$. We map the evolution of a massive Dirac band in the vicinity of the Fermi level as a function of the magnetic field. We find that electronic states at different lattice momenta exhibit different Zeeman energy shifts, giving rise to an anomalous $g$ factor that peaks around the Dirac point. Our work provides a momentum-resolved visualization of Dirac band curvature manipulated by a magnetic field, which will be relevant to other topological kagome magnets.

The Zeeman effect, first observed as the atomic spectral line splitting under an applied magnetic field, has found applications across different branches of physics. The effect is rooted in the linear coupling between the total magnetic moment and the external magnetic field $B$ detectable via the field-induced energy shift $-g (\mu_B|B|) J \cdot B$, where $J$ is the spin–orbit angular moment of the atomic states and $g$ is a dimensionless number called the Landé $g$ factor. Bloch electrons in quantum solids also exhibit Zeeman coupling to the magnetic field, set by an effective $g$ factor that takes on a more complex functional form. In general, the $g$ factor can vary with the overall band structure, spin–orbit coupling, carrier concentration, and exchange and correlation effects$^{3,4}$, and it is fully described by a matrix that is a function of the crystal momentum $k$ (ref. 9). Although the theoretical framework behind the $g$ factor momentum dependence was established around the middle of the last century$^{1-4}$, variations of the $g$ factor with momentum in ordinary metals and semiconductors are largely negligible.

It is theorized, however, that certain materials exhibit a strongly momentum-dependent effective $g$ factor, which can lead to substantial field-induced modifications to the electronic band structure. In graphene and topological materials, such as twisted bilayer graphene, topological and Chern insulators, and Dirac and Weyl semimetals, the $g$ factor can, in principle, be enhanced by the nontrivial topological band dispersion through momentum-dependent orbital magnetizations$^{14-16}$. In magnetic and/or spin-textured systems, the external magnetic field may further give rise to unusual Zeeman energy shifts and strong renormalization of the bands, enabling spin-direction-driven field-sensitive tuning of the electronic band structure$^{16-17}$ and topological phase transitions$^{18}$. However, a direct momentum-resolved insight into the field-driven changes to the electronic band structure by experimentally measuring the shift of the quantum states in an external magnetic field has been difficult to achieve. In this Letter we report the experimental discovery that the effective electron $g$ factor associated with a massive Dirac band, determined directly by measuring the Zeeman energy shift in a kagome magnet YMn$_6$Sn$_6$, is strongly momentum-dependent.

Using spectroscopic-imaging scanning tunnelling microscopy (STM), we map a massive Dirac dispersion crossing the Fermi level, as a function of magnetic field. We find that the energy of each state within the Dirac band shifts linearly with applied field. By charting the field-induced shift of each electronic state within the band, we map a strong $k$-dependent renormalization of the band curvature due to Zeeman coupling, which can be explained by a $k$-dependent electron $g$ factor.

Layered crystalline materials composed of atoms arranged on a lattice of corner-sharing triangles (kagome lattice) have emerged as a rich material platform to study electronic correlations and nontrivial topology$^{14-18}$. A simple tight-binding calculation of an isolated kagome plane, taking into account nearest-neighbour hopping, reveals that the system hosts Dirac points near the edge of the Brillouin zone and a dispersionless flat band$^{15,16}$. Motivated by theoretical predictions of new electronic states, such as a fractional quantum Hall state without external field$^{17-19}$, spin liquid phases$^{20}$, tunable Weyl nodes$^{21}$ and a Wigner crystal$^{22}$, exotic phenomena have been intensely sought after and realized in several kagome compounds$^{22-28}$.

YMn$_6$Sn$_6$ is a layered material from a rare-earth (Re) family of kagome magnets, ReMn$_6$Sn$_6$, with an in-plane hexagonal structure $(a=b=0.554\text{ nm}, c=0.901\text{ nm})^{29}$. It goes through an antiferromagnetic transition at $T_N \approx 350\text{ K}$ (ref. 29), which doubles the unit cell shown in Fig. 1a along the $c$ axis. As the temperature decreases further, the spin structure of YMn$_6$Sn$_6$ becomes helical$^{28-32}$. A key ingredient of the crystal structure is the kagome lattice of Mn atoms, characterized by a network of corner-sharing Mn triangles, sandwiched between alternating stacks of Sn$^-$–Sn$^-$–Sn$^-$ and YSn layers (Fig. 1a). We cleaved and investigated the surface of YMn$_6$Sn$_6$ single crystals grown by the flux method as described in ref. 29 (Methods). STM topographs typically show surface terminations parallel to the $a$–$b$ plane (Fig. 1c), although we occasionally found small patches.

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of other crystalline facets as well (Extended Data Fig. 1). We classify the a–b-plane termination layers based on the morphology of the STM topographs, topographic terrace heights and average differential conductance (dV/dI) spectra (Fig. 1d–i, Extended Data Figs. 2a,b and 3 and Methods). These experimental identifications were further verified by simulated STM images from first-principles calculations (Extended Data Fig. 2).

Here we focus on the electronic structure of the Mn kagome surface. The average dV/dI spectra show a sharp peak near the Fermi level at εF ≈ −7.0 ± 0.1 meV (Figs. 1g and 2). This peak is spatially homogeneous, which rules out localized bound states tied to crystal defects (Supplementary Fig. 1). To gain further insight, we performed spectroscopic measurements with magnetic field B applied parallel to the c axis (perpendicular to the sample surface). We tracked an identical 40-nm-square field of view at different fields and compared the average dV/dI spectra (Fig. 2a,b). Other than the spectral peak at εF, we observed no additional peaks that could be interpreted as Landau levels up to the highest field used in our experiments (Supplementary Fig. 4). If the peak at εF was associated with a doubly spin-degenerate band, we would expect the two spin branches to shift in opposite directions due to spin Zeeman coupling; correspondingly, the spectral peak in dV/dI would split. In contrast to this, we find that the peak does not split (nor broaden) up to the maximum applied B of 8 T (Fig. 2a,b). This points towards a singly spin-polarized band. For a fixed spin orientation regardless of external B, we would expect a diverging shift of εF for a field applied parallel and antiparallel to the c axis. This possibility is again

Fig. 2 | Magnetic field dependence of the spectral peak in the density of states. a, Average dV/dI spectra acquired in an identical region of the Mn layer from −8 T (antiparallel to the c axis) to 8 T (parallel to the c axis). The green and grey arrows indicate the approximate positions of the dip and the peak at εF, respectively. b, dV/dI spectra from a in a different colour scale for visualization purposes. White circles in b denote peak positions extracted by fitting (Supplementary Fig. 2). c, Scatterplot of εF as a function of applied magnetic field, showing linear dispersion with slope α = 0.31 ± 0.03 meV T\(^{-1}\). STM set-up conditions: I\(_{\text{set}}\) = 70 pA, V\(_{\text{sample}}\) = 30 mV, V\(_{\text{set}}\) = 1 mV.

Fig. 1 | Crystal structure and different surface terminations. a, Schematic of the YMn\(_2\)Sn\(_4\) crystal structure. b, A topographic line profile across consecutive terraces, extracted along the red line in the inset. Each step shows a single unit cell step height of about 0.87 nm. c, Different possible surface terminations associated with the a–b plane: Mn, YSn\(^3\), Sn\(^2\) and Sn\(^1\) layers. d–i, STM topographs (d–f) and average dV/dI spectra (g–i) on the kagome Mn layer (d,g). Sn\(^2\) layer (e,h) and Sn\(^1\) layer (f,i). A spectral peak labelled as εF is present near the Fermi level on the Mn layer, but not on the Sn\(^2\) or Sn\(^1\) layers, which show a small ‘dip’ at the Fermi level as often observed on non-kagome terminations in related systems\(^{23,34}\). Dashed lines in g–i denote the approximate energy range of the linearly dispersing differential conductance. STM set-up conditions: I\(_{\text{set}}\) = 10 pA, V\(_{\text{sample}}\) = 1 V (b); I\(_{\text{set}}\) = 70 pA, V\(_{\text{sample}}\) = 30 mV (d–f); I\(_{\text{set}}\) = 300 pA, V\(_{\text{sample}}\) = 500 mV, V\(_{\text{set}}\) = 2 mV (g–i).
contradictory to our observations, as $E_p$ shifts in the same direction, even if the direction of $B$ is reversed (Fig. 2). Such a shift is consistent with a minority spin state in an electronic state polarized by the external field itself. Conventional Zeeman spin coupling characterized by electron $g$ factor $g_e \approx 2$ would dictate that the state shifts to lower energy with velocity $1/2 g_e \mu_B \approx 0.058 \text{meV}\,\text{T}^{-1}$ ($\mu_B$ is one Bohr magneton). In contrast, we find that the state moves to a higher energy by fivefold the velocity of $0.31 \text{meV}\,\text{T}^{-1}$ (Fig. 2c). This amounts to a magnetic moment of $-5.34 \mu_B$ or a large effective $g$ factor of $g_p \approx -10.7$. Note that the dip spectral feature (green arrows in Fig. 2a) shifts to higher energies faster than the adjacent peak $E_p$, which indicates a complex response of the electronic structure to the magnetic field in this system, reaching beyond a rigid band shift.

A more detailed momentum-space picture of the low-energy electronic band structure can be obtained from spectroscopic-imaging STM (SI-STM). This method, also commonly referred to as quasi-particle interference (QPI) imaging, relies on the elastic scattering and interference of electrons, detectable as static modulations in $dI/dV(r, V)$ maps. Let us first focus on the SI-STM measurement at zero magnetic field. Real-space $dI/dV(r, V)$ maps show intense spatial modulations (Fig. 3b, top row) with a nearly isotropic Fourier space wavevector $q_1$ that rapidly evolves with energy (Fig. 3b, bottom row). The inset in $b$ in the FT at $-10 \text{mV}$ shows the radially averaged linecut (red points indicate the data) with clearly distinguishable scattering peaks (the two blue curves are Gaussian fits). $c$, Radially averaged linecut in the FTs of the $dI/dV$ maps. Label $q_1$ in $b$ and $c$ denotes the prominent scattering vector in momentum-transfer space, which shows a nearly isotropic spatial signature and a rapid evolution with energy. The dispersion velocity of the band associated with the QPI signal in $c$ is calculated to be $373.6 \pm 9.4 \text{meV}\,\text{Å}$. The centre pixel of the FT linecut in $c$ has been artificially suppressed to emphasize other features. $d$, Schematic of the band structure based on ARPES, with the Dirac point and the saddle point indicated by arrows. The Fermi level was shifted down by approximately 30–40 meV compared to the ARPES measurement to match the energy of the STM spectral features. STM set-up conditions: $I_{\text{set}} = 70 \text{pA}, V_{\text{sample}} = 30 \text{mV}, V_{\text{exc}} = 2 \text{mV}, B = 0 \text{T}$ (a,b); $I_{\text{set}} = 140 \text{pA}, V_{\text{sample}} = 100 \text{mV}, V_{\text{exc}} = 4 \text{mV}, B = 0 \text{T}$ (c).
row), disappearing around \(-15\) meV (Fig. 3c and Supplementary Fig. 3). From the QPI data we extract the corresponding dispersion of the band crossing the Fermi level (Fig. 3c). Note that dispersion shows a linear energy–momentum dependence at higher energies, but deviates from the linear dependence near the band bottom and acquires substantial curvature (outlined by the dashed line in Fig. 3c). By comparing the Fermi velocity and the morphology of the QPI signal to the low-energy electronic band structure near the Fermi level measured recently by angle-resolved photoemission spectroscopy (ARPES)\(^{33}\) (Fig. 3d), we can attribute the QPI dispersion to gapped Dirac cones at \(K\), where the gap is a consequence of spin–orbit coupling. The sharp density-of-states peak at \(E_g\) in Fig. 2 is then probably associated with the van Hove saddle point at \(M\) of a different band, located only a few millielectronvolts above the Dirac band bottom (Figs. 3d and 4a,b). The partial flat band observed by ARPES in the second Brillouin zone is located \(-20\) meV below the Dirac point\(^{33}\), and it is thus unlikely to be associated with the sharp spectral feature at \(E_g\). We note that the lower half of the Dirac cone cannot be unambiguously discerned in our data, possibly due to a decreased quasiparticle lifetime away from the Fermi energy. This is similar to the markedly weaker electronic signature of the upper half of the Dirac cone in the related material TbMn\(_6\)Sn\(_6\) (ref. \(^{34}\)), which is also positioned further away from zero energy.

Next we applied an external magnetic field of up to 8 T in both the \(c\) and \(\overline{c}\) directions. Similar to the behaviour of the spectral peak at \(E_g\) (Fig. 2), the band dispersion did not split and shifted to higher energies independent of the direction of the magnetic field, thus pointing towards a magnetization-polarized origin of a minority band (Fig. 4a–d). The quantitative field-induced energy shift of the individual electronic states within the band at different momenta \(k\) surrounding the \(K\) point can be plotted as a function of the magnetic field (Fig. 4e–g). We stress that the shift of all states shows a linear dependence following the Zeeman effect, but with \(k\)-dependent slopes. The states near the band bottom (\(k=0\) referenced to the \(K\) point) exhibit the most pronounced evolution with energy (Fig. 4e). As the momentum moves away from the Dirac band bottom, the energy shift velocity gradually decreases, but saturates at a constant value of \(0.48\) meV T\(^{-1}\) (Fig. 4e–h). This momentum-dependent
band shift is remarkable and has not been observed experimentally before. It directly reveals a \( k \)-dependent \( g \) factor for Bloch electrons. Writing the energy shift as \( \Delta E(k) = -\frac{i}{\hbar} g \mu_B B \), where \( B \) is the magnitude of the magnetic field, we extract the \( k \)-dependent \( g \) factor \( g_k \) and plot it as a function of \( k \) in Fig. 4. Clearly, the behaviour of the \( g \) factor follows \( g_k = g_0 + \delta g_k \), describing a constant background contribution of \( |g_k| \approx 15 \) and a \( k \)-dependent Lorentzian-like \( \delta g_k \) that ranges from 0 to 12, peaking around the K point at \( K=0 \).

Theoretically, there are several contributions that can, in principle, lead to the \( k \)-dependent shift of Dirac states that is most pronounced near the bottom of the observed electron-like part of the massive Dirac bands. First, note that the Dirac mass gap is present already at zero magnetic field. Similarly to the momentum-dependent \( g \) factor in certain graphene-based structures\(^{25-28} \), nontrivial Berry curvature associated with the massive Dirac fermions generates an orbital magnetic moment, and the orbital Zeeman coupling to the magnetic field along the \( c \) axis is consistent with the observed energy shift linear in the magnetic field (Zeeman shift). Our calculation in a simplified model of this system, based on a single-orbital, spin-polarized kagome lattice model, confirms the large orbital magnetic moments, primarily localized near the Dirac points at \( K \) where the Berry curvature concentrates (Supplementary Discussion 4 and Supplementary Fig. 6c). The coupling to the applied magnetic field can account for the linear shift in the energy of the states resulting from the orbital Zeeman effect, in a \( k \)-dependent fashion set by \( k \)-dependent orbital moments. Second, as the magnetization measurements of \( \text{YMn}_6\text{Sn}_6 \) indicate canting of spins towards the field direction (Supplementary Fig. 7), the Kane–Mele-type spin–orbit coupling due to the in-plane orbital motion can naturally enhance the Dirac gap as the magnetization gains an out-of-plane component (Supplementary Discussion 5)\(^{29} \). It is thus expected that the observed rounding of the Dirac band bottom (Fig. 4c) contains a contribution from the spin–orbit coupling as the out-of-plane magnetic field is applied. Finally, we note that magnetic exchange coupling, as well as the intricate spin textures\(^{30,31} \) and its evolutions under a \( c \)-axis field (currently under intensive investigation), can, in principle, also lead to additional field-induced band renormalization\(^{25} \). Future experiments probing the evolution of electronic properties under a vector magnetic field may provide the further insights that are necessary to disentangle the different underlying contributions to the observed momentum-dependent \( g \) factor and fully understand the magnetic field tuning of the Dirac states revealed here.

Transport measurements have reported unconventional responses to the magnetic field in a variety of bulk single crystals\(^{25-28} \), but a direct insight into the field-induced electronic band structure change has been challenging. Our experiments reveal a momentum-dependent effective \( g \) factor associated with a massive Dirac band in a kagome magnet. We note that, although previous experiments reported a rigid field-induced shift of the spectral peaks associated with flat portions of the kagome bands\(^{34,35} \), the bottom of the Dirac band\(^{32} \) or localized states\(^{41} \), our work also reveals continuous changes in the associated electronic momenta. Although the origin of the Dirac mass at zero field is beyond the scope of this work, it is conceivable that the spin–orbit coupling in \( \text{YMn}_6\text{Sn}_6 \) stems from beyond the Kane–Mele scenario and requires the full crystal symmetry and both in-plane and out-of-plane spin components. In addition, the spin Berry phase due to non-coplanar spin textures may also contribute to the zero-field Dirac mass, as in the case of \( \text{Fe}_3\text{Sn}_2 \) (ref. 31). Owing to the close proximity of the Dirac points to the Fermi level in \( \text{YMn}_6\text{Sn}_6 \), the anomalous physics reported here could potentially be probed by transport experiments. Given the complexity of the bulk electronic and magnetic structures of this system, which is composed of at least eight kagome layers within a magnetic unit cell\(^{31} \), a complete quantitative understanding of the momentum-dependent \( g \) factor will probably require a thorough many-body theoretical calculation, taking into account the helical magnetic structure and contributions from other Dirac and flat bands, including exchange effects\(^{46} \). As the \( k \)-dependent \( g \) factor can produce a large band renormalization, it in turn enables fine tunability of the electronic structure by an external magnetic field, including the possibility of driving the many-body ground states through different topological phase transitions. Because the magnetic order at the surface of \( \text{YMn}_6\text{Sn}_6 \) has broken time-reversal symmetry, the gapped Dirac band should carry a non-zero Chern number, pointing to \( \text{YMn}_6\text{Sn}_6 \) being a topological kagome lattice Chern magnet.

**Online content**

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Methods

Sample growth and characterizations. Single crystals of YMn6Sn6 were grown using Sn flux. Yttrium lumps (purity 99.99%), Mn granules (purity 99.9%) and Sn grains (purity 99.99%) with a molar ratio of Y: Mn: Sn = 1: 6: 30 were placed in an almost saturated Sn atmosphere. The sealed quartz ampoule was heated to 1,273 K and held for 24 h, then cooled slowly to 873 K at a rate of 5 K h\(^{-1}\). The ampoule was then taken out of the furnace and decanted with a centrifuge to separate YMn6Sn6 crystals from excess Sn flux.

STM experiments. Single crystals of YMn6Sn6 were cleaved in ultrahigh vacuum and immediately inserted into the STM head, where they were held at ~4.5 K during measurements. Of the more than 40 crystals in which cleaving was attempted, we successfully cleaved and approached seven different samples. Two of these showed large enough Mn terraces for high-resolution spectroscopic mapping (Fig. 1 and Supplementary Fig. 8). STM data were acquired using a customized Unisoku USM1300 STM system. Spectroscopic measurements were made using a standard lock-in technique with 915-Hz frequency and bias excitation as detailed in the figure captions. The zero energy of our measurements was checked, and no noticeable difference was observed as a function of magnetic field (Supplementary Fig. 5). The used STM tips were home-made chemically etched tungsten tips, annealed in UHV to a bright orange colour before STM experiments.

Surface layer identification and additional surface terminations observed. To corroborate the nature of the cleavage planes in Fig. 1 we considered the atomic structure in STM topographs, step heights between consecutive terraces, average \(dI/dV\) spectra and theoretical simulations (Supplementary Discussion 1). For the \(a-b\)-plane surface termination, consecutive terraces typically showed a single unit cell step height along the \(c\) direction (Fig. 1b) and a hexagonal lattice structure (Fig. 1c and Extended Data Fig. 3). The Mn kagome plane and Sn\(^2\) termination both show a honeycomb-like pattern (Fig. 1d,e), but markedly different average \(dI/dV\) spectra (Fig. 1g,h), and Sn\(^1\) termination is characterized by a triangular lattice of atomic features (Fig. 1f). We identify the Mn kagome layer (Fig. 1d) from the half-unit-cell step height between consecutive terraces we sometimes encounter, which, electronically and topographically, look nearly identical (Extended Data Fig. 2). We further note that the Mn kagome plane imaged here looks qualitatively similar to the kagome planes of several related materials, such as Co\(_n\)Sn\(_m\)S\(_n\) (ref. 32), Fe\(_2\)Sn\(_3\) (ref. 33), Mn\(_2\)Sn\(_3\) (ref. 34) and TbMn\(_6\)Sn\(_6\) (ref. 34). Although comparable to the topograph of the Mn layer, the STM topograph of the Sn\(^2\) plane shows clearly resolved individual atoms within each hexagon (Fig. 1c) and looks tantalizingly similar to the honeycomb Sn layer of Fe\(_2\)Sn\(_3\) (ref. 32) and the single-layer stanene grown on copper\(^{43}\). Finally, we identify the Sn\(^1\) termination (Fig. 1f) based on the atomic structure of the layer and the step height offset from the Mn layer (Extended Data Fig. 2). These experimental identifications are consistent with simulated STM images of the corresponding terminations from our first-principles electronic structure calculations (Extended Data Fig. 2).

STM topographs of the cleaved surfaces of YMn\(_6\)Sn\(_6\) exhibit a variety of other surface morphologies beyond the \(a-b\)-surface terminations shown in Fig. 1. Several examples of such morphologies are shown in Extended Data Fig. 1; these all exhibit unidirectional features and are extremely slanted. By fitting each raw STM topograph with a second-degree polynomial, we can determine the slopes of each of these planes, which range from 20 to 30°. For comparison, the inclination angles of the \(a-b\)-plane surface terminations shown in the main text, determined by the same fitting procedure, are less than 2°. We thus hypothesize that these slanted surfaces, although atomically ordered, are not \(a-b\)-plane surface terminations, but are other, possibly reconstructed, cleavage planes.

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

Code availability
The computer code used for data analysis is available upon request from the corresponding author.

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Author contributions
H.Li and H.Z. performed STM experiments. H.Li analysed the STM data. Q.W. and Q.Y. synthesized and characterized the samples under the supervision of H.Lei, K.J. and Z.W. performed orbital magnetization calculations. N.-N.Z. and K.L. performed simulations of the STM topographs. H.Li, H.Z., Z.W. and I.Z. wrote the manuscript, with input from all authors. I.Z. supervised the project.

Competing interests
The authors declare no competing interests.

Additional information

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Extended Data Fig. 1 | Other surface terminations observed in YMn$_6$Sn$_6$. (a–d) STM topographs of other crystalline facets observed on surfaces with large slopes (approximate surface angle is shown in the upper left). The upper right corner in each panel shows the Fourier transform of the corresponding topograph, with the most prominent peaks circled in red. For comparison, all STM topographs of the $ab$-plane in the main text show nominal angle of approximately 2 degrees or less, a typical uncertainty of gluing the sample to the holder by silver epoxy. STM setup condition: (a, b) $I_{\text{set}} = 10$ pA, $V_{\text{sample}} = 400$ mV; (c) $I_{\text{set}} = 70$ pA, $V_{\text{sample}} = 30$ mV; (d) $I_{\text{set}} = 10$ pA, $V_{\text{sample}} = 400$ mV.
Extended Data Fig. 2 | Surface identification based on step heights and theoretical simulations. (a) STM topograph of consecutive steps. (b) Topographic line profile taken along the red line denoted in (a). The total height between the bottom layer and the top layer is 8.7 Å, which is a unit cell height. Based on the step heights and the nature of surface morphologies over each layer, surface terminations are identified as Mn, Sn, Mn and Mn (the tallest to the shortest terrace). (c–e) Theoretical simulations of STM topographs of Mn, Sn1 and Sn2 terminations at 30 mV bias. (f–h) Experimental STM topographs of Mn, Sn1 and Sn2 surface terminations. STM setup condition: (c–e) simulated $V_{\text{sample}} = 30$ mV. (f–h) $I_{\text{set}} = 70$ pA, $V_{\text{sample}} = 30$ mV.
Extended Data Fig. 3 | Bias-dependent STM topographs. STM topographs of (a-d) Mn and (e-h) Sn’ surface terminations as a function of bias. The lower right corner of (a,e) shows a Fourier transform of the topograph in that panel, with hexagonal Bragg peaks circled in red. STM topographs measured at different bias over the same region of the sample show qualitatively the same surface morphology, regardless of the bias. STM setup condition: (a) $I_{\text{set}} = 100$ pA, $V_{\text{sample}} = 1$ V. (b) $I_{\text{set}} = 100$ pA, $V_{\text{sample}} = 50$ mV. (c) $I_{\text{set}} = 100$ pA, $V_{\text{sample}} = -50$ mV. (d) $I_{\text{set}} = 100$ pA, $V_{\text{sample}} = -1$ V. (e) $I_{\text{set}} = 30$ pA, $V_{\text{sample}} = 1$ V. (f) $I_{\text{set}} = 30$ pA, $V_{\text{sample}} = 50$ mV. (g) $I_{\text{set}} = 30$ pA, $V_{\text{sample}} = -50$ mV. (h) $I_{\text{set}} = 30$ pA, $V_{\text{sample}} = -1$ V.