Electronic nematicity without charge density waves in titanium-based kagome metal

Layered crystalline materials that consist of transition metal atoms on a kagome network have emerged as a versatile platform for the study of unusual electronic phenomena. For example, in the vanadium-based kagome superconductors $\text{A}_2\text{V}_3\text{Sb}_5$ (where $\text{A}$ can stand for $\text{K}$, $\text{Cs}$ or $\text{Rb}$), there is a parent charge density wave phase that appears to simultaneously break both the translational and rotational symmetries of the lattice. Here we show a contrasting situation, where electronic nematic order—the breaking of rotational symmetry without the breaking of translational symmetry—can occur without a corresponding charge density wave. We use spectroscopic-imaging scanning tunnelling microscopy to study the kagome metal $\text{CsTi}_3\text{Bi}_5$ that is isostructural to $\text{A}_2\text{V}_3\text{Sb}_5$ but with a titanium atom kagome network. $\text{CsTi}_3\text{Bi}_5$ does not exhibit any detectable charge density wave state, but a comparison to density functional theory calculations reveals substantial electronic correlation effects at low energies. In comparing the amplitudes of scattering wave vectors along different directions, we discover an electronic anisotropy that breaks the sixfold symmetry of the lattice, arising from both in-plane and out-of-plane titanium-derived $d$ orbitals. Our work uncovers the role of electronic orbitals in $\text{CsTi}_3\text{Bi}_5$, suggestive of a hexagonal analogue of the nematic bond order in Fe-based superconductors.
occurs concomitant with the formation of the CDW. It has also been reported that rotational symmetry breaking is a CDW-driven phenomenon. Theoretically, rotational symmetry breaking in \( AV_3Sb_5 \) could be explained by the interlayer stacking order of the CDW phase, which is intrinsically three-dimensional and in turn selects one out of the three lattice directions as the preferred one. While other distinct experimental signatures of rotational symmetry breaking, such as charge stripe order, anisotropy in elasto-transport measurements and unidirectional coherent quasiparticles, are found to onset at even lower temperatures, they all inevitably occur within the system, where both rotational and translational symmetries appear to be already broken at CDW onset (refs. 34, 35). As such, within the liquid crystal classification of electronic phases, the parent electronic order of \( AV_3Sb_5 \) could be regarded as a ‘smectic’ phase. This brings the following fundamental questions: To what extent is rotational symmetry breaking in kagome metals tied to the CDW order itself? Can it exist in the absence of translational symmetry breaking to form a genuine electronic ‘nematic’ phase?

Recently, a new family of Ti-based kagome metals \( ATi_3Bi_x \) (where \( A = Cs \) or Rb) has been synthesized in the same crystalline structure as \( AV_3Sb_5 \), but with a kagome net of Ti atoms replacing the V and Bi substituting for Sb. While some studies suggested that \( CsTi_3Bi_x \) may be superconducting below critical temperature \( T_c = 4 \) K (ref. 36), similarly to \( CsV_3Sb_5 \), superconductivity has not been confirmed by other studies yet (Extended Data Fig. 1). From the fermiology perspective, \( ATi_3Bi_x \) also hosts van Hove singularities at M points, but these are now positioned well above the Fermi level unlike the equivalent features in \( AV_3Sb_5 \) that appear in closer proximity to the Fermi level. In contrast to \( AV_3Sb_5 \), experiments on \( ATi_3Bi_x \) revealed no obvious anomalies in magnetization, resistivity and heat capacity that could be attributed to a CDW state, which has been widely observed in \( AV_3Sb_5 \).

This in turn provides an opportunity to explore new electronic phenomena in a non-magnetic kagome metal in the absence of translational symmetry breaking CDW. Our scanning tunnelling microscopy (STM) measurements of cleaved \( CsTi_3Bi_x \) bulk single crystals are consistent with the absence of a CDW in this system. Moreover, we use spectroscopic-imaging scanning tunnelling microscopy (SI-STM) to discover a pronounced rotational symmetry breaking in the electronic signal. In particular, we observe substantial directionality in the amplitudes of scattering wave vectors between the three nominally identical high-symmetry directions. We find that anisotropic spectral weight associated with different portions of the bands can naturally explain unidirectionality in our SI-STM measurements. In sharp contrast to \( AV_3Sb_5 \), where rotational symmetry breaking signatures occur within the CDW phase that already breaks the rotational symmetry, the reduction of the rotational symmetry here clearly occurs in the absence of the concomitant translational symmetry breaking, thus forming a pure electronic nematic phase. This is akin to the electronic nematic phase in Fe-based high-temperature superconductors but on a hexagonal lattice.

Bulk single crystals of \( CsTi_3Bi_x \) exhibit a hexagonal crystalline structure \((a = 5.7 \text{ Å}; c = 9.2 \text{ Å})\) composed of Ti–Bi slabs stacked between alkali Cs layers (Fig. 1a,b). The Ti kagome net is interlaced with Bi atoms and sandwiched between two other hexagonal Bi layers (Fig. 1a). Due to the air-sensitive nature of the samples, we handle sample preparation in an inert environment of an argon glove box and minimize subsequent exposure to air during the transfer to the STM chamber (Methods). The samples are cleaved in ultra-high vacuum at cryogenic temperatures and immediately inserted into the STM for measurements. Similarly to the cleaving structure extensively explored in \( AV_3Sb_5 \), \( CsTi_3Bi_x \) crystals should naturally cleave between the Cs and Bi layers, resulting in two different surface terminations. Indeed, STM

---

**Fig. 1** Crystalline structure and topographic characterization of \( CsTi_3Bi_x \). (a) A three-dimensional ball model of the crystalline structure of \( CsTi_3Bi_x \). (b) A two-dimensional ball model showing relative positions of different atoms in the a–b plane. (c) A 20 × 20 nm² STM topograph of \( CsTi_3Bi_x \) taken at 4.5 K showing Cs islands on top of a complete Bi termination. (d) A topographic line profile taken along the red dashed line in c across a step between the two terminations.

---

| Image 1 | Image 2 | Image 3 | Image 4 | Image 5 |
|---------|---------|---------|---------|---------|
| ![Image 1](image1.png) | ![Image 2](image2.png) | ![Image 3](image3.png) | ![Image 4](image4.png) | ![Image 5](image5.png) |

---

**Article** https://doi.org/10.1038/s41567-023-02176-3
topographs reveal the two main types of atomically ordered surface morphologies: an incomplete layer that we attribute to the Cs layer on top of a complete honeycomb-like surface structure, which we ascribe to the Bi layer (Fig. 1c,d). While STM topographs of both surfaces can show atomic resolution (Fig. 1e,f), we note that the Cs surface tends to form disconnected island-like patches, and that large complete regions of the Cs surface are difficult to locate. In contrast, Bi-terminated layers are robust, and we can find large, flat areas for imaging (Fig. 2a).

Fourier transforms (FTs) of low-temperature STM topographs of the Bi surface show a hexagonal lattice with wave vectors $\mathbf{a}_{\text{STM}} = b_{\text{STM}} = 5.8$ Å (inset in Fig. 2a), which are consistent with the expected bulk lattice constants. Aside from atomic Bragg peaks $Q_{\text{bragg}}$, no additional peaks are detected along high-symmetry directions in the FTs of STM topographs (Fig. 2b). This is consistent with the absence of a CDW in CsTi$_3$Bi$_5$, in contrast to the CDW widely observed in the STM topographs of the $\text{AV}_2\text{Sb}_3$ family acquired over a wide range of biases.  

---

**Fig. 2 | Large-scale electronic band structure and the absence of CDW in CsTi$_3$Bi$_5$.**

a. STM topograph taken over a clean Bi-terminated region. The bottom right inset is its FT showing a hexagonal lattice. b. FT linecuts starting from the centre of the FT in the inset in a along the three atomic Bragg peak directions, showing the absence of the 2 × 2 CDW peaks present in cousin CsV$_3$Sb$_5$. c,d. An example of a normalized conductance map $G(r, V) = (dI/dV)/I(r, V)$ taken approximately over the region in a (c) and its FT (d). Yellow, orange and red arrows denote the three dominant scattering wave vectors $\mathbf{q}_1$, $\mathbf{q}_2$, and $\mathbf{q}_3$, respectively. e. Radially averaged linecut extracted from the FTs of $G(r, V)$ maps. The data near 0 mV are artificially suppressed due to the divergent behaviour of $G(r, V)$ for near-zero bias and d, DFT-calculated band structure of CsTi$_3$Bi$_5$ along high-symmetry directions $K$–$\Gamma$–$K$–$M$–$\Gamma$, with different colours denoting either Bi (red) or Ti (blue) orbital character. g. An enlarged FT linecut near the Fermi level taken along the purple dashed line in d across the three wave vectors.

h. Comparisons between dispersions of $\mathbf{q}_i$ (where $i = 1, 2, 3$) and equivalent vectors extracted from the DFT band structure for the momentum along $z$ axis $k_z = 0$ (blue) and $k_z = 0.5$ (green lines). STM data points are determined by fitting a Gaussian peak function to the linecut in g at energies separated by 10 meV increments. STM setup conditions: $I_{\text{set}} = 200$ pA and $V_{\text{sample}} = 100$ mV (a); $I_{\text{set}} = 600$ pA, $V_{\text{sample}} = 100$ mV and $V_{\text{exc}} = 4$ mV (c, d, g); $I_{\text{set}} = 1$ nA, $V_{\text{sample}} = 500$ mV and $V_{\text{exc}} = 10$ mV (e). Data were taken at 4.5 K.
Fig. 3 | Twofold symmetric electronic signature. a, FT of a normalized differential conductance \( I(r, V = -14 \text{ mV}) \) map taken on the Bi surface. Atomic Bragg peaks are circled in black. Blue and red double arrows depict where \( \mathbf{q}_1 \) and \( \mathbf{q}_2 \) wave vector amplitudes are extracted. b,c, Angle-dependent FT amplitudes of \( \mathbf{q}_1 \) (b) and \( \mathbf{q}_2 \) (c) plotted in polar coordinates on the same scale (see also Extended Data Fig. 9 for an estimate of spectral weight anisotropy). The purple dashed lines in a–c denote the apparent \( C_2 \)-symmetric axis. The results of the atomic Bragg peak analysis can be found in Extended Data Fig. 10. d, Representative energy-dependent FTs also showing the twofold symmetric scattering signature. e,f, Momentum-transfer space positions of the wave vectors \( \mathbf{q}_1 \) (e) and \( \mathbf{q}_2 \) (f) along the three \( \Gamma - K \) directions. As can be seen, both \( \mathbf{q}_1 \) and \( \mathbf{q}_2 \) are slightly larger along direction \( A \) compared to the other two directions \( B \) and \( C \), which appear nearly indistinguishable. Directions \( A, B \) and \( C \) are defined in a. Error bars in e and f represent the standard errors obtained by Lorentzian fits to the data. g,h, FTs of the \( I(r, V) \) maps taken over two different regions on the same \( \text{CsTi}_3\text{Bi}_5 \) sample 2, denoted as regions I (g) and II (h), both of the Bi surface. Region II is found after a manual rotation of the whole sample by \( 120^\circ \), while the tip wire in both cases remains the same. Atomic Bragg peaks are circled in black. i,j, Angle-dependent FT amplitudes of \( \mathbf{q}_1 \) (blue) and \( \mathbf{q}_2 \) (red line) plotted in polar coordinates for data in g (i) and h (j). The symmetry axis is approximately along the \( 15^\circ \) angle in i and along the \( 25^\circ \) angle in j. STM setup conditions: \( I_{\text{set}} = 300 \text{ pA}, V_{\text{sample}} = 30 \text{ mV} \) and \( V_{\text{exc}} = -4 \text{ mV} \) (a,d); \( I_{\text{set}} = 300 \text{ pA}, V_{\text{sample}} = 30 \text{ mV} \) and \( V_{\text{exc}} = -4 \text{ mV} \) (e,f); \( I_{\text{set}} = 200 \text{ pA}, V_{\text{sample}} = -10 \text{ mV} \) and \( V_{\text{exc}} = 2 \text{ mV} \) (g); \( I_{\text{set}} = 500 \text{ pA}, V_{\text{sample}} = 10 \text{ mV} \) and \( V_{\text{exc}} = 5 \text{ mV} \) (h).

To gain insight into the electronic band structure of \( \text{CsTi}_3\text{Bi}_5 \), we perform SI-STM measurements, which are rooted in elastic scattering and interference of electrons detectable as periodic modulations in differential conductance \( dI/dV(r, V) \) maps. Quantitative assessment of different scattering wave vectors can be more easily obtained from discrete FTs of \( dI/dV(r, V) \) maps, which provide momentum-space information. The FTs of \( dI/dV(r, V) \) maps acquired at high energies on the Bi surface of \( \text{CsTi}_3\text{Bi}_5 \) show a series of dispersive scattering wave vectors (Fig. 2c,d). To understand their origin, we focus on the three concentric wave vectors that appear over a wide energy range, labelled \( \mathbf{q}_1, \mathbf{q}_2 \) and \( \mathbf{q}_3 \) (Fig. 2d,e,g). To visualize the dispersion of these vectors, we first extract radially averaged FT linecuts of normalized \( dI/dV \) maps as a function of energy (Fig. 2e). We also extract a linear FT linecut to determine the Fermi velocities of different bands (Fig. 2f). We find that they exhibit large and positive dispersion velocities: \( 5,200 \pm 230 \text{ mV Å} \) for \( \mathbf{q}_1, 3,330 \pm 70 \text{ mV Å} \) for \( \mathbf{q}_2 \), and \( 2,370 \pm 100 \text{ mV Å} \) for \( \mathbf{q}_3 \) (Extended Data Fig. 2). We compare the morphology and the dispersion of these wave vectors in SI-STM data to density functional theory (DFT) calculations of \( \text{CsTi}_3\text{Bi}_5 \), including spin–orbit coupling (Fig. 2f and Methods). DFT calculations show several electron-like bands centred within the Bi-derived \( p \)-orbital pocket at \( \Gamma \) due to a close agreement in dispersion velocities, Fermi wave vectors and the band bottom energies between theory and experiment (Fig. 2e–h). The outer wave vectors, \( \mathbf{q}_1 \) and \( \mathbf{q}_2 \), arise from scattering between \( Ti\ d_{xy}/d_{yz} \) band and Bi \( p \) band (\( \mathbf{q}_1 \)) and intraband scattering within the Ti \( d_{xy}/d_{yz} \) band (\( \mathbf{q}_2 \)). We can conclude this based on the morphology of the wave vectors and comparable dispersion velocities (Fig. 2g,h). Interestingly, however, there is a sizeable difference in the energy of Ti-derived bands compared to what is expected from theory (Fig. 2h), suggesting a relative energy offset between the Bi- and Ti-derived bands not captured by theory. We note that this cannot be explained by a rigid band shift of all bands induced by small surface doping. While an unusual band-selective surface doping could contribute to this behaviour, we deem this to be unlikely, as relevant portions of bands involved in scattering in our STM data appear to be pushed down by electron doping.
in a comparable manner. This in turn may already hint at additional Coulomb interaction and correlation effects in this material.

A visual inspection of the FTs leads us to another intriguing observation—while each wave vector should appear to be the same in a comparable manner. This in turn may already hint at additional Coulomb interaction and correlation effects in this material.

To quantify the angle-dependent amplitude variation, we plot the FT of \( \mathbf{q} \) along one direction (purple dashed line) appear to be more intense compared to the other two nominally equivalent directions. In particular, \( \mathbf{q}_a \) and \( \mathbf{q}_b \) along one direction (purple dashed line) appear to be more intense compared to the other two nominally equivalent directions. To quantify the angle-dependent amplitude variation, we plot the FT amplitudes of the peaks as a function of angle (Fig. 3b,c). Both show the same dominant axis. To confirm the robustness of this observation, we evaluate potential tip anisotropy artefacts that could cause the apparent FT intensity anisotropy. We first note that we repeated the experiment on multiple samples with different tip wires (Fig. 3a,g,h and Extended Data Fig. 3). We then show that the intensity anisotropy is robust to rotating the scan direction by 90° (Extended Data Fig. 4). We stress that we have changed the microscopic arrangement of atoms at the tip apex in situ by fast scanning and bias pulsing multiple times, and we still have consistently observed the same symmetry axis over the same area of the sample (Extended Data Fig. 3). In addition to the intensity anisotropy of the scattering wave vectors, we also find angular dependence of the scattering wave vector magnitudes (Fig. 3e,f). While this angular variation of vector magnitudes is only at the order of a few percent, it directly demonstrates small anisotropy of the associated constant energy contour (CEC), elongated along the same direction as the symmetry axis observed in vector intensity plots in Fig. 3a. We note that we have not been able to locate a domain wall across which the dominant axis would rotate by 120°, possibly due to the the apparent FT intensity anisotropy.

To confirm the robustness of this observation, we proceed to map the apparent symmetry axis before and after a manual rotation of the sample by 120° while keeping the same orientation of the tip wire (Fig. 3g–j). If breaking of the rotational symmetry observed in Fig. 3 was an artefact due to the tip anisotropy, we would expect to still find angular dependence of the scattering wave vector magnitudes and its auto-correlations (bottom panels). Different thicknesses of lines in each CEC represent a varying spectral weight of the associated portion of each band.
observe the same orientation of a tip artefact-induced C\textsubscript{3} symmetric axis in equivalent SI-STM scans. Importantly, however, the orientation of the symmetry axis is now also rotated by 120° (Fig. 3g compared to Fig. 3h), pointing against a spurious C\textsubscript{3} symmetric signal caused by tip anisotropy. Taken together, our experiments revealed that the rotational symmetry of the electronic structure of CsTi\textsubscript{3}Bi\textsubscript{5} is reduced from C\textsubscript{3} symmetry to C\textsubscript{2} symmetry.

We proceed to examine the orbital structure of electronic states associated with scattering wave vectors in more detail by examining CECs obtained from our DFT calculations. The Fermi surface consists of several closed contours (inset in Fig. 4b), which show a weak k\textsubscript{f} dependence (Extended Data Fig. 5). The small circular pocket centred at Γ is primarily composed of Bi\textsubscript{p} orbitals, while the rest of the bands are largely derived from different Ti d orbitals. For pedagogical purposes, we illustrate what subsets of bands in the CEC play a role in the emergence of each set of features seen in the experimental FT (top panels in Fig. 4c–f), and then we perform auto-correlation of the CEC to mimic the features in the experimental FT (bottom row in Fig. 4c–f).

For example, as shown in Fig. 4c, we find that the previously discussed wave vectors in Fig. 2 can be captured by scattering between p\textsubscript{z} orbitals (q\textsubscript{z}), p\textsubscript{x} and d\textsubscript{xy} orbitals (q\textsubscript{xy}), and d\textsubscript{yz} orbitals (q\textsubscript{yz}). In addition to these previously discussed features, FTs at low energies also show a series of elongated quasi-one-dimensional features near the Brillouin zone edge (black rectangle in Fig. 4a). These features clearly disperse with energy (Fig. 4h,i) and can be explained by an interplay of scattering and interference of electrons primarily from d\textsubscript{xy}/d\textsubscript{yz} orbitals (Fig. 4d–f). These elongated scattering wave vectors near the Brillouin zone edge appear more intense along one direction (black rectangle in Fig. 4a) and look tantalizingly similar to the quasi-one-dimensional vectors in CsV\textsubscript{3}Sb\textsubscript{5}, (refs. 6,11,35) that also arise from d-orbital scattering. Putting together the aforementioned scattering processes (Fig. 4b,g) beautifully reproduces the fine structure of the experimental FT (Fig. 4a,i). Importantly, to capture the observation that scattering appears markedly unidirectional, a simple model taking into account a variation of the quasiparticle spectral weight associated with different bands is sufficient (Fig. 4c–f).

Rotational symmetry breaking in V-based kagome superconductors has been reported to be CDW-driven\textsuperscript{30} and intimately tied to the smectic CDW phase onset\textsuperscript{13,25}, which could in principle be explained by the interlayer stacking order of the three-dimensional CDW phase\textsuperscript{20,21}. In contrast to the smectic order in AV\textsubscript{3}Sb\textsubscript{5}, our experiments revealed the emergence of a pure electronic nematic state in CsTi\textsubscript{3}Bi\textsubscript{5}, in the absence of the CDW phase in cousin system CsV\textsubscript{3}Sb\textsubscript{5}. As such, our work shifts the paradigm of rotational symmetry breaking in kagome metals beyond any ties to CDW stacking and charge stripe orders. Magnetotransport experiments further support the underlying rotational symmetry breaking in CsTi\textsubscript{3}Bi\textsubscript{5} (ref. 46). While no obvious anomalies are detected in zero-field resistivity, magnetization or heat capacity\textsuperscript{5}, it is possible that the signal may be too weak to be detected in such experiments.

There are several aspects of electronic nematicity in CsTi\textsubscript{3}Bi\textsubscript{5} that stand out compared to other families of materials. First, in most other systems where evidence for electronic nematicity is observed, such as Cu- and Fe-based superconductors, strong electronic correlation and proximity to magnetism make it difficult to determine whether the nematic order is driven by spin or orbital degrees of freedom. The kagome metals have paramagnetic spin susceptibility without any magnetism, either localized or itinerant. It is therefore surprising that nematicity arises in CsTi\textsubscript{3}Bi\textsubscript{5} which offers a fresh opportunity to explore an orbital-driven nematic order. Second, electronic nematic order in hexagonal systems is a new subject with a potential for realizing three-state Potts nematic order\textsuperscript{47}. The commonly studied materials, with the exception of the twisted Moiré structures\textsuperscript{38–41}, do not offer this possibility. However, in contrast to nematicity in Moiré heterostructures where rotational symmetry breaking occurs upon a partial filling of localized electrons occupying flat bands, electronic nematicity in CsTi\textsubscript{3}Bi\textsubscript{5} clearly affects delocalized Ti-derived d orbitals. It is interesting to note that angular anisotropy of q\textsubscript{z}, connecting p\textsubscript{z} orbitals only, appears much smaller compared to q\textsubscript{xy} and q\textsubscript{yz}, which involve d orbitals (Extended Data Fig. 6), further supporting the notion of electronic nematicity from d electronic orbitals. How correlations of the electronic orbitals can lead to Potts nematicity is a new frontier of quantum materials, and our work provides ground work in that direction.

It is possible that the smectic CDW in CsV\textsubscript{3}Sb\textsubscript{5} and the nematic order in CsTi\textsubscript{3}Bi\textsubscript{5} share similar mechanisms, but the emergence of CDWs in CsV\textsubscript{3}Sb\textsubscript{5} occurs due to stronger effective interactions, as correlation effects in CsV\textsubscript{3}Sb\textsubscript{5} may be enhanced due to higher density of states in proximity to van Hove filling\textsuperscript{25,30}. However, we note that the two systems show fundamental differences, including the differing strengths of spin–orbit coupling and different energies of van Hove singularities.

In contrast to CsV\textsubscript{3}Sb\textsubscript{5}, that has a pronounced lattice instability contributing to the emergence of the smectic CDW phase, no such instability exists in the phonon calculations of CsTi\textsubscript{3}Bi\textsubscript{5} (Extended Data Fig. 7). This points towards electronically driven nematic state in CsTi\textsubscript{3}Bi\textsubscript{5}. A natural explanation could lie in orbital ordering, such as ferro-orbital or bond orbital ordering, where extended Coulomb interactions play a significant role\textsuperscript{5}. The multi-orbital nature of CsTi\textsubscript{3}Bi\textsubscript{5} with multiple Ti d bands crossing the Fermi level is akin to that of Fe-based superconductors, where several Fe d bands span the Fermi level. It is conceivable that the electronic nematicity uncovered here could be described by the nematic bond order proposed to explain the electronic nematic phase in FeSe\textsuperscript{42,43} but on a hexagonal kagome lattice distinct from the nearly square lattice of FeSe. Interestingly, if similarities between normal states of FeSe and CsTi\textsubscript{3}Bi\textsubscript{5} extend to the superconducting regime, it would be of high interest to explore in future experiments if and how electronic nematicity in the normal state affects the mechanism of Cooper pairing in this family of kagome metals.

**Online content**

Any methods, additional references, Nature Portfolio reporting summaries, source data, extended data, supplementary information, acknowledgements, peer review information; details of author contributions and competing interests; and statements of data and code availability are available at https://doi.org/10.1038/s41567-023-02176-3.

**References**

1. Ortiz, B. R. et al. New kagome prototype materials: discovery of KV\textsubscript{3}Sb\textsubscript{5}, RbV\textsubscript{3}Sb\textsubscript{5}, and CsV\textsubscript{3}Sb\textsubscript{5}. *Phys. Rev. Mater.* **3**, 094407 (2019).
2. Ortiz, B. R. et al. CsV\textsubscript{3}Sb\textsubscript{5}: a Z\textsubscript{2} topological kagome metal with a superconducting ground state. *Phys. Rev. Lett.* **125**, 247002 (2020).
3. Yin, Q. et al. Superconductivity and normal-state properties of kagome metal RbV\textsubscript{3}Sb\textsubscript{5} single crystals. *Chin. Phys. Lett.* **38**, 037403 (2021).
4. Ortiz, B. R. et al. Superconductivity in the Z\textsubscript{2} kagome metal KV\textsubscript{3}Sb\textsubscript{5}. *Phys. Rev. Mater.* **5**, 034801 (2021).
5. Jiang, Y.-X. et al. Unconventional chiral charge order in kagome superconductor KV\textsubscript{3}Sb\textsubscript{5}. *Nat. Mater.* **20**, 1353–1357 (2021).
6. Zhao, H. et al. Cascade of correlated electron states in the kagome superconductor CsV\textsubscript{3}Sb\textsubscript{5}. *Nature* **599**, 216–221 (2021).
7. Li, H. et al. Rotation symmetry breaking in the normal state of a kagome superconductor KV\textsubscript{3}Sb\textsubscript{5}. *Nat. Phys.* **18**, 265–270 (2022).
8. Guo, C. et al. Switchable chiral transport in charge-ordered kagome metal CsV₃Sb₅. Nature 611, 461–466 (2022).
9. Wu, Q. et al. Simultaneous formation of two-fold rotation symmetry with charge order in the kagome superconductor CsV₃Sb₅ by optical polarization rotation measurement. Phys. Rev. B 106, 205109 (2022).
10. Mielke, C. et al. Time-reversal symmetry-breaking charge order in a kagome superconductor. Nature 602, 245–250 (2022).
11. Chen, H. et al. Roton pair density wave in a strong-coupling kagome superconductor. Nature 599, 222–228 (2021).
12. Liang, Z. et al. Three-dimensional charge density wave and surface-dependent vortex-core states in a kagome superconductor CsV₃Sb₅. Phys. Rev. X 11, 031026 (2021).
13. Li, H. et al. Observation of unconventional charge density wave without acoustic phonon anomaly in kagome superconductors AV₃Sb₅ (A = Rb, Cs). Phys. Rev. X 11, 031050 (2021).
14. Kang, M. et al. Twofold van Hove singularity and origin of charge order in topological kagome superconductor CsV₃Sb₅. Nat. Phys. 18, 301–308 (2022).
15. Xiang, Y. et al. Twofold symmetry of c-axis resistivity in topological kagome superconductor CsV₃Sb₅ with in-plane rotating magnetic field. Nat. Commun. 12, 6727 (2021).
16. Yang, S.-Y. et al. Giant, unconventional anomalous Hall effect in the metallic frustrated magnet candidate, KV₃Sb₅. Sci. Adv. 6, eabb6003 (2020).
17. Xu, H.-S. et al. Multiband superconductivity with sign-preserving order parameter in kagome superconductor CsV₃Sb₅. Phys. Rev. Lett. 127, 187004 (2021).
18. Uykur, E., Ortiz, B. R., Wilson, S. D., Dressel, M. & Tsirlin, A. A. Optical detection of the density-wave instability in the kagome metal KV₃Sb₅. NPJ Quantum Mater. 7, 16 (2022).
19. Feng, X., Jiang, K., Wang, Z. & Hu, J. Chiral flux phase in the kagome superconductor AV₃Sb₅. Sci. Bull. 66, 1384–1388 (2021).
20. Tan, H., Liu, Y., Wang, Z. & Yan, B. Charge density waves and electronic properties of superconducting kagome metals. Phys. Rev. Lett. 127, 046401 (2021).
21. Denner, M. M., Thomale, R. & Neupert, T. Analysis of charge order in the kagome metal AV₃Sb₅ (A = K, Rb, Cs). Phys. Rev. Lett. 127, 217601 (2021).
22. Lin, Y.-P. & Nandkishore, R. M. Complex charge density waves at Van Hove singularity on hexagonal lattices: Haldane-model phase diagram and potential realization in the kagome metals AV₃Sb₅ (A = K, Rb, Cs). Phys. Rev. B 104, 045122 (2021).
23. Park, T., Ye, M. & Balents, L. Electronic instabilities of kagome metals: saddle points and Landau theory. Phys. Rev. B 104, 035142 (2021).
24. Christensen, M. H., Biro, T., Andersen, B. M. & Fernandez, R. M. Theory of the charge density wave in AV₃Sb₅ kagome metals. Phys. Rev. B 104, 214513 (2021).
25. Zhou, S. & Wang, Z. Chern Fermi pocket, topological pair density wave, and charge-4e and charge-6e superconductivity in kagome superconductors. Nat. Commun. 13, 7288 (2022).
26. Christensen, M. H., Biro, T., Andersen, B. M. & Fernandez, R. M. Loop currents in AV₃Sb₅ kagome metals: multipolar and toroidal magnetic orders. Phys. Rev. B 106, 144504 (2022).
27. Liu, Z. Z. et al. Charge-density-wave-induced bands renormalization and energy gaps in a kagome superconductor RbV₃Sb₅. Sci. Rev. X 11, 041010 (2021).
28. Hu, Y. et al. Topological surface states and flat bands in the kagome superconductor CsV₃Sb₅. Sci. Bull. 67, 495–500 (2022).
29. Luo, H. et al. Electronic nature of charge density wave and electron-phonon coupling in kagome superconductor KV₃Sb₅. Nat. Commun. 13, 273 (2022).
Methods
Bulk single-crystal synthesis and characterization
Bulk single crystals of CsBi₃Ti₅ have been synthesized, as described in ref. 37.

STM experiments
Bulk single crystals of CsTi₃Bi₅ were transported from the University of California, Santa Barbara, to Boston College in a sealed glass vial filled with argon gas to prevent air degradation. The vial was opened in an argon glove box at Boston College, where the sample was glued to a UNISOKU-style STM sample holder, and a cleave bar was then attached to the sample. We quickly transferred the prepared sample from the glove box to an ultra-high vacuum lock load within seconds and cold-cleaved it before putting it into the microscope. STM data were acquired using a custom UNISOKU USM300 microscope at about 4.5 K. Spectroscopic measurements were performed using a standard lock-in technique with 910 Hz frequency and bias excitation noted in the figure captions. The STM tips used were homemade chemically etched tungsten tips, annealed in ultra-high vacuum to a bright orange colour before the STM experiments. To remove the effects of small piezoelectric and thermal drifts during the acquisition of data, we applied the Lawler–Fujita drift-correction algorithm on all our data, which aligns the atomic Bragg peaks in STM topographs to be equal in magnitude and oriented 60° apart.

Ruling out STM tip artefacts
To rule out the possibility of tip artefacts artificially inducing the electronic anisotropy, we characterized the electronic symmetry axis on the Bi surface, retracted the tip and manually rotated the sample 120°. After the tip re-approaches on the sample, we took equivalent dI/dV maps again. As shown, the electronic symmetry axis also rotated by 120° (Fig. 3g,h). Furthermore, four different Bi regions were explored on two different CsTi₃Bi₅ samples, and all of them showed the C₂-symmetric scattering and interference patterns (several examples shown in Extended Data Fig. 3). In these regions, we manually performed bias pulsing of the tip on the Bi surface (which changes the microscopic arrangement of atoms at the tip apex) to change the tip, and we examined the data before and after the tip change. We have demonstrated that similar anisotropic scattering and interference patterns can be resolved in the same region even with different tips (Extended Data Fig. 3). As such, we deem that the C₂-symmetric electronic signature cannot be explained by the STM tip artefacts.

DFT calculation
The structure of CsTi₃Bi₅ was fully relaxed with the Vienna Ab initio Simulation Package (VASP), where the Perdew–Burke–Ernzerhof (PBE)-type generalized gradient approximation was employed to mimic the electron–electron interaction. The cutoff energy for the plane-wave basis set was 300 eV, and a k mesh of 12 × 12 × 6 was used to sample the Brillouin zone. The DFT-D3 van der Waals correction was considered in the relaxation. After that, the Full-Potential Local-Orbital (FPLO) software was employed to calculate the band structure under PBE approximation with spin–orbit coupling considered. A k-mesh size of 12 × 12 × 6 was used. The Fermi surface of CsTi₃Bi₅ was calculated from the tight-binding Hamiltonian, as obtained from the FPLO Wannier fitting. The Ti d and Bi p orbitals were used as the Wannier basis set.

Data availability
The data supporting the findings of this study are available from Zenodo at https://doi.org/10.5281/zenodo.8092076 and also upon request from the corresponding author. Source data are provided with this paper.

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

References
55. Kresse, G. & Furthmüller, J. Efficiency of ab-initio total energy calculations for metals and semiconductors using a plane-wave basis set. Comput. Mater. Sci. 6, 15–50 (1996).
56. Perdew, J., Burke, K. & Ernzerhof, M. Generalized gradient approximation made simple. Phys. Rev. Lett. 77, 3865–3868 (1996).
57. Grimme, S., Antony, J., Ehrlich, S. & Krieg, H. A consistent and accurate ab initio parametrization of density functional dispersion correction (DFT-D) for the 94 elements H-Pu. J. Chem. Phys. 132, 154104 (2010).
58. Koepernik, K. & Eschrig, H. Full-potential nonorthogonal local-orbital minimum-basis band-structure scheme. Phys. Rev. B 59, 1743–1757 (1999).

Acknowledgements
I.Z. gratefully acknowledges the support from the National Science Foundation (NSF), Division of Materials Research 2216080. S.D.W. and B.R.O. acknowledge financial support from the US Department of Energy (DOE), Office of Basic Energy Sciences, Division of Materials Sciences and Engineering under Grant No. DE-SC0020305. This work used facilities supported via the University of California, Santa Barbara, NSF Quantum Foundry funded via the Quantum Materials Science, Engineering and Information program under award DMR-1906325. Z.W. acknowledges the support of the US Department of Energy, Basic Energy Sciences Grant No. DE-FG02-99ER45747 and the Cottrell Singular Exceptional Endeavors of Discovery Award No. 27856 from Research Corporation for Science Advancement. D.W. and D.J. acknowledge the support from the Bavaria California Technology Center Grant 7 [2021-2]. B.Y. acknowledges the financial support by the European Research Council (ERC Consolidator Grant 'NonlinearTopo', No. 815869) and the ISF - Personal Research Grant (No. 2932/21).

Author contributions
STM experiments and data analysis were performed by H.L., with the help from S.C. and B.R.O. D.W. synthesized and characterized the samples under the supervision of S.D.W. and D.J. Z.W. and K.Z. provided theoretical input on data interpretation. H.T. and B.Y. performed band structure calculations. H.L., S.D.W., Z.W. and I.Z. wrote the paper, with input from all authors. I.Z. supervised the project.

Competing interests
The authors declare no competing interests.

Additional information
Extended data is available for this paper at https://doi.org/10.1038/s41567-023-02176-3.

Supplementary information The online version contains supplementary material available at https://doi.org/10.1038/s41567-023-02176-3.

Correspondence and requests for materials should be addressed to Ilija Zeljkovic.

Peer review information Nature Physics thanks Youichi Yamakawa and the other anonymous reviewer(s) for their contribution to the peer review of this work.

Reprints and permissions information is available at www.nature.com/reprints.
Extended Data Fig. 1 | Low-temperature differential conductance spectra. Average $dI/dV$ spectra on the Bi surface acquired at approximately 350 mK (left) and 2 K (right) temperature showing the absence of a clear superconducting gap. STM setup conditions: (left) $I_{\text{set}} = 500$ pA, $V_{\text{sample}} = 5$ mV, $V_{\text{exc}} = 0.1$ mV, 0 T; (right) $I_{\text{set}} = 1$ nA, $V_{\text{sample}} = 5$ mV, $V_{\text{exc}} = 0.1$ mV, 0 T.
Extended Data Fig. 2 | Comparison of QPI wave vectors and theoretical band structure. (a) Fourier transform (FT) of a Bi surface $dI/dV(r, V = 20 \text{ mV})$. Three pink arrows denote $q_1$, $q_2$, and $q_3$ scattering wave vectors. (b) FT linecut taken along the blue dashed arrow in (a). The three dispersive QPI wave vectors are clearly seen, which are denoted as $q_1$, $q_2$, and $q_3$. (c) The summary of energy dependent vector position in $k$ space. Data in (b,c) was converted from momentum-transfer $q$-space to $k$-space by $1/2$ multiplication factor. Each point in (c) is extracted by fitting a Gaussian peak function to the linecuts in (b). The error bars represent standard errors of peak positions. The comparison of the linearly fitted band velocities with electronic band velocities in (d) points to the origin of different $q$-vectors as discussed in the main text: $q_1$ comes from intrascattering within the Bi orbital derived $\Gamma$ pocket, $q_3$ arises from the intrascattering within the Ti orbital derived, and $q_2$ is likely due to scattering between the Bi and the Ti orbital derived bands. (d) The DFT calculated band structure along $\Gamma$–$K$ direction. STM setup conditions: $I_{\text{set}} = 600 \text{ pA}$, $V_{\text{sample}} = 100 \text{ mV}$, $V_{\text{exc}} = 4 \text{ mV}$ (a).
Extended Data Fig. 3 | Robustness of observed anisotropy to using different tips. (a) Fourier transforms (FTs) of $dI/dV$ maps taken on the same Bi surface region with different tips (left and right panel). Blue and red arrows denote $q_2$ and $q_3$ wave vectors. (b) Angle-dependent polar plots of $q_2$ and $q_3$ amplitudes before and after the intentional tip change. Red line corresponds to the $q_3$ amplitude and the blue line corresponds to the $q_2$ amplitude. (c,d) Equivalent to (a,b) but for a different region of the sample. It is evident that amplitudes exhibits a strong $C_2$ symmetric signal on different regions with different tips, with the dominant axis pointing along a $\Gamma$-$K$ direction in each dataset. Note that the scan frame was slightly rotated between region A and B, which is why the lattice also appears rotated between panels (a) and (c). These experiments provide further evidence that $C_2$ symmetric electron scattering is not an artifact caused by the tip shape. STM setup condition: $I_{set} = 200$ pA, $V_{sample} = -10$ mV, $V_{exc} = 2$ mV (region A); $I_{set} = 500$ pA, $V_{sample} = 10$ mV, $V_{exc} = 5$ mV (region B).
Extended Data Fig. 4 | Robustness of observed anisotropy to scan direction.
(a) A 35 nm square STM topograph taken over the region A of sample 2.
(b, c) Fourier transforms (FTs) of $\frac{dI}{dV}$ maps taken over the region in (a) scanned
along (b) vertical and (c) horizontal direction at -7 mV bias. Atomic Bragg peaks
are circled in black. Blue and red arrows denote $q_2$ and $q_3$.
(d, e) The amplitudes of $q_2$ and $q_3$ plotted as a function of angle in polar axis corresponding to data in
(b, c). It is clear that both patterns are $C_2$-symmetric, and the symmetry axis
does not change due to a different scanning direction. STM setup condition:
$I_{\text{set}} = 200$ pA, $V_{\text{sample}} = -10$ mV, $V_{\text{exc}} = 2$ mV (a-c).
Extended Data Fig. 5 | Theoretically calculated constant energy contours. Constant energy contours at -50 meV (left), 0 mV (middle) and +50 meV (right) calculated by DFT for $k_z = 0$ (red lines) and $k_z = 0.5$ (blue lines).
Extended Data Fig. 6 | Additional scattering wave vector dispersions. (a) Fourier transform of a normalized differential conductance $L(V = -14\,\text{mV})$ map taken on the Bi surface (also shown in Fig. 3a of the main text). Atomic Bragg peaks are circled in black. (b) Momentum-transfer space positions of the wave vectors $q_1$, $q_2$, and $q_3$ along the three $\Gamma$-$K$ directions. Error bars in (b) represent the standard errors obtained by Lorentzian fits to the data. The anisotropy of $q_1$ is much smaller compared to the other two wave vectors.
Extended Data Fig. 7 | Calculated phonon dispersions. Calculated phonon dispersion in (a) CsTi$_3$Bi$_5$ and (b) CsV$_3$Sb$_5$ using the same procedure as described in Tan et al, *PRL* 127, 046401 (2021). Imaginary (negative) phonon frequencies in (b) marked by red arrows correspond to the breathing modes of the kagome lattice, which likely contributes to the CDW distortions in CsV$_3$Sb$_5$. Corresponding instability in CsTi$_3$Bi$_5$ is however absent.
Extended Data Fig. 8 | STM topograph spanning both surface terminations. (a) STM topograph spanning both the Cs and the Bi surface termination. (b) Zoom-in on a smaller region showing the atomic structure of both surfaces across a step edge. The lattice of small white circles shows the positions of individual Cs atoms based on the known atomic structure of the Bi termination. It can be seen that Cs atoms on the Cs termination reside on top of bright features in the topograph. STM setup conditions: $I_{\text{set}} = 200$ pA, $V_{\text{sample}} = 100$ mV.
Extended Data Fig. 9 | Relating QPI anisotropy and spectral weight anisotropy. (a) Form factor function $f(\Theta, N = 1.5) = 1 + (\cos(2\Theta) - 1)/(2N)$ used to simulate the spectral weight anisotropy. (b) Autocorrelation of the constant energy contour (CEC) at zero energy, which is multiplied by $f(\Theta, N = 1.5)$ prior to the autocorrelation procedure. This qualitatively simulates the $C_2$-symmetric QPI by adding an angle-dependent function to the weight of the Fermi surfaces to break their $C_6$ symmetry. $\Theta$ is angle calculated with respect to the $\Gamma$-$K$ direction. As seen in (c), to approximately get the ratio of $q_3$ wave vectors along the two directions to be 2, $N$ factor is about 1.5. In this scenario, the spectral weight varies by about a factor of 2 as seen in panel (a), consistent with QPI anisotropy.
Extended Data Fig. 10  | Energy-dependence of the signal at atomic Bragg peak wave vectors. (a) Fourier transform of the L-map from the main text Fig. 3 showing the three atomic Bragg peak positions used for the analysis. The peak labeled A is along a mirror symmetry axis. (b) Amplitude of the atomic Bragg peaks as a function of energy from dI/dV maps. (c,d) Kagome lattice simulation. We generate a Kagome lattice real space density of states (DOS) distribution by using a simple trigonometric function: 

\[
|\psi(r)|^2 = \sum_{j=A,B,C} |\psi(j)|^2 \prod_{i=0,1,2} (1 + \cos \frac{4\pi}{\sqrt{3}} e_i \cdot [\frac{2\pi(1-\alpha)}{6} \cdot (r-r_j)]),
\]

where \(r_A=(0,0); r_B=(1,0)\) a/2, \(r_C=(1,\sqrt{3})\) a/4. \(e_i = (1,0)\), \(R[\theta]\) is the rotation matrix, and \(a\) is the lattice constant. When \(|\psi(A)|^2 = |\psi(B)|^2 = |\psi(C)|^2\), the DOS at the three sub-lattices are all the same. However, when (c) \(|\psi(A)|^2 > |\psi(B)|^2 = |\psi(C)|^2\) (A site has a higher DOS than B and C), or (d) \(|\psi(A)|^2 < |\psi(B)|^2 = |\psi(C)|^2\) (A site has a smaller DOS compared to B and C), the simulated lattices break the \(C_6\) symmetry. By tuning \(|\psi(B)\|^2/|\psi(A)\|^2\) from 1 to 0, sites B and C gradually fade away, and the lattice becomes triangular, as shown in (e). By tuning \(|\psi(A)\|^2/|\psi(B) = C)\|^2\) from 1 to 0, the A site gradually disappears as shown in (f), which ultimately leads to a rectangular lattice. Fourier transforms (FTs) of the simulated topographs are shown in the second row of (e,f). In both cases, the lattice symmetry varies due to the variation of A site DOS relative to that of B and C sites. (g,h) Atomic Bragg peak \(Q_{\text{Bragg}}\) intensities depending on the DOS ratios between the sub-lattices.