An anomaly in quantum field theory is defined as the breaking of a classically allowed symmetry when quantum effects are turned on. The first example, called the Adler–Bell–Jackiw or chiral anomaly, was discovered in 1968 in the process of understanding the observed rapid decay rate of the neutral pion into two photons. When fermions are massless, they segregate into independent right- and left-handed groups (the light up and down quarks in pions are treated as massless). However, once the massless fermions are coupled to electromagnetic fields, the chiral symmetry is violated, along with the appearance of an axial current (see the Supplementary Information). Calculations of the axial current accurately matched the decay rate of the pion (refs 4,5).

Subsequently, anomalies have appeared in diverse phenomena occurring at vastly different energy scales. In 1983, it was proposed that the chiral anomaly should be observable in crystals as a large, negative longitudinal magnetoresistance (LMR)6.

In the past two decades, research on Dirac electrons in semimetals has grown rapidly into a major field of activity, especially in topological quantum matter. The realization of 3D Dirac states in semimetals has led to a surge of experiments investigating the LMR. Five years ago, the discovery of Dirac and Weyl semimetals led to many experiments investigating this phenomenon. In this Review, we critically assess LMR experiments in the Dirac–Weyl semimetals Na3Bi, GdPtBi, ZrTe5, Cd3As2, and TaAs, which have shown signatures of the chiral anomaly, and discuss possible current-jetting artefacts. The focus is on Dirac and Weyl nodes that are rigorously symmetry protected. Other experiments, such as non-local transport, thermopower, thermal conductivity and optical pump–probe response, are also reviewed. Looking ahead, we anticipate what can be gleaned from improved LMR experiments and new experiments on the thermal conductivity and optical response. An expanded purview of the chiral anomaly is provided in the Supplementary Information.

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In the past two decades, research on Dirac electrons in semimetals has grown rapidly into a major field of activity, especially in topological quantum matter. The realization of 3D Dirac states in semimetals has led to a surge of experiments investigating the LMR and revived intense interest in the chiral anomaly and its observability in crystals (refs 6). The findings of negative LMR fall into two distinct groups. In the semimetals Na3Bi and GdPtBi, which have low carrier mobilities, the LMR is very large (the longitudinal resistivity can decrease by a factor of 10) and monotonic, settling down to a B-independent constant in high magnetic field B. In the second group, comprised of the Weyl semimetals TaAs and NbAs, which have very high mobilities, the LMR is strongly non-monotonic; it is weakly negative (0.1–2%) in small B but becomes positive above 1 T. Inconsistencies soon raised strong concerns that the observed LMR in the second group is an artefact caused by field-induced inhomogeneous current flow known as current jetting. A test to distinguish intrinsic LMR from current-jetting artefacts confirmed that the negative LMR is intrinsic in the first group but artefactual in the second.

In this Review, we critically assess LMR experiments in the Dirac and Weyl semimetals Na3Bi, GdPtBi, ZrTe5, Cd3As2, and TaAs, which have shown signatures of the chiral anomaly, and discuss possible current-jetting artefacts in LMR data. Alternative approaches to investigating the chiral anomaly, including experiments on non-local transport, thermopower, thermal conductivity and optical pump–probe response, are also reviewed. In the Supplementary Information, we provide a brief discussion of the chiral anomaly in the broader context of high-energy physics and relativistic quantum field theory, as well as the anomaly's starring role at the nexus of quantum physics and differential geometry. For a complementary review on Dirac semimetals, we direct readers to ref 21, and for a review of the chiral magnetic effect (CME) in heavy-ion collisions, see ref 22.

Symmetry protection and Dirac–Weyl nodes
We start by discussing chiral symmetry. In general, if the Lagrangian \( \mathcal{L} \) of a system is invariant under a symmetry operation, we can derive a conservation law that...
Key points

- The chiral anomaly describes the conversion of left-moving massless fermions to right-moving ones in the presence of electromagnetic fields and has long been predicted to be observable in crystals.
- The chiral anomaly was predicted in 1983 to be observable in crystals and the discovery of Dirac and Weyl semimetals with symmetry-protected nodes now enables tests of the prediction.
- Longitudinal magnetoresistance measurements can provide evidence of the chiral anomaly but can also contain artefacts caused by current jetting.
- Other experimental measurements of the chiral anomaly include thermal conductivity and optical response.
- These experiments join a long line of anomaly experiments that extend over a vast range of energy scales in many subfields of physics.

relates a conserved charge to the associated current (Noether’s theorem). An example is the invariance of \( \mathcal{L} \) under a global (uniform) rotation of the phase of the field \( \Psi \to \Psi e^{i\theta} \). The conservation law resulting from this global gauge invariance is also familiar. The rate of change of the local charge density \( \rho \) equals the (negative) divergence of the charge-current density \( \mathbf{J} \). Our interest is in systems that harbour two fields, which we call \( \Psi_\uparrow \) (L for left-moving) and \( \Psi_\downarrow \) (R for right-moving). These fields may describe electron populations \( N_\uparrow \) and \( N_\downarrow \) occupying 1D states that disperse strictly to the left and to the right, respectively. In most cases, \( \mathcal{L} \) is invariant under a phase rotation of both \( \Psi_\uparrow \) and \( \Psi_\downarrow \) by the same angle \( \theta \). The resulting conservation law relates the total charge \( N_\uparrow + N_\downarrow \) to the total current density \( \mathbf{J} = \mathbf{J}_\uparrow + \mathbf{J}_\downarrow \), as mentioned.

When the states are massless, a second symmetry arises. The Lagrangian is unchanged if we rotate the phase of \( \Psi_\uparrow \) and \( \Psi_\downarrow \) in opposite directions, such that \( \Psi_\uparrow \to \Psi_\uparrow e^{i\theta} \) and \( \Psi_\downarrow \to \Psi_\downarrow e^{-i\theta} \). Invariance under this (axial) rotation is called chiral or axial symmetry. The conservation law derived from chiral symmetry relates the population difference \( N_\uparrow - N_\downarrow \equiv N^\perp \) to the axial current density \( \mathbf{J}^\perp = \mathbf{J}_\uparrow - \mathbf{J}_\downarrow \). For details, see the Supplementary Information. It is perhaps unsurprising that the population difference is conserved. Massless fermions split into two independent populations \( N_\uparrow \) and \( N_\downarrow \) that are separately conserved. However, once we couple the system to the electromagnetic field, both the chiral symmetry and the conservation of \( N^\perp \) are violated. In the 1D example, this just expresses charge pumping in an electric field \( E \) from the uphill-moving branch (the branch climbing the potential gradient) to the downhill branch. The former, say, \( N_\uparrow \), decreases at the rate \( \dot{N}_\uparrow = -(e/2\pi)E \), whereas the latter increases as \( \dot{N}_\downarrow = (e/2\pi)E \). Instead of being conserved, the difference current \( \mathbf{J}^\perp \) now acquires a source term \( (e/\pi)E \), which is the net pump rate. The breaking of chiral symmetry is known as the chiral anomaly and the source term the anomaly term.

Our emphasis in this Review is on bulk Dirac and Weyl nodes that are symmetry protected, just like the nodes in graphene (an elementary discussion of protection is given in the Supplementary Information). Early searches for Dirac nodes in crystals were based on tuning a single parameter (usually pressure or alloy content \( x \))\(^1\). These nodes are, however, unprotected by symmetries and, hence, unstable — they stay ungapped at only a single value of the tuning parameter. Although a weak negative LMR observed in \( \mathrm{Bi}_1 - \mathrm{Sb}_x \) (\( x \approx 0.03 \)) was interpreted as the chiral anomaly\(^2\), the current understanding is that unprotected Dirac nodes do not lead to protected Weyl nodes\(^3-6\). Subsequent investigations have focussed exclusively on materials in which the Dirac and Weyl nodes are symmetry protected. In the Dirac semimetal \( \mathrm{Na}_3\mathrm{Bi} \), band inversion leads to two band crossings lying on the symmetry axis \( \Gamma - A (|| k_z) \) of the point group element \( \mathcal{C}_3 \). As discussed below, the \( \mathcal{C}_3 \) symmetry, together with time-reversal invariance (TRI) and inversion symmetry, protect the Dirac nodes against gap formation. We call these stable nodes. In the presence of a finite Zeeman field (TRI broken), each Dirac node separates into two daughter Weyl nodes with opposite chirality. Crucially, because the Weyl nodes inherit the symmetry protection against gap formation, the chiral nature of their lowest Landau level provides the setting for the chiral anomaly to appear.

The search for 3D protected Dirac states in semimetals\(^7-11\) initially focussed on the two symmetries, TRI and inversion symmetry. These are sufficient to protect 3D Dirac nodes that are located at the time-reversal invariant momenta (TRIM) \( \mathbf{K} \) (which satisfy \( \mathbf{K} = -\mathbf{K} + \mathbf{G} \), with \( \mathbf{G} \) a reciprocal lattice vector). The predicted materials, notably \( \beta \)-crystobalite \( \mathrm{BiO}_2 \)(Ref.\(^11\)), are chemically unstable. Subsequently, it was realized that including point-group symmetry confers protection of Dirac nodes anywhere along a symmetry axis (e.g. the rotation axis of \( \mathcal{C}_n \) with \( n = 3, 4, 6 \)). Freeing the node from a TRIM led to the prediction that the two semimetals \( \mathrm{Na}_3\mathrm{Bi} \)(Ref.\(^12\)) and \( \mathrm{Cd}_3\mathrm{As}_2 \)(Ref.\(^13\)) are Dirac semimetals, with nodes at wavevectors \( \mathbf{k}^\perp_D \) protected by symmetry under \( \mathcal{C}_n \) and \( \mathcal{C}_l \) rotations, respectively. In the case of \( \mathrm{Na}_3\mathrm{Bi} \), the workhorse for the chiral anomaly, it has been shown that, with the choice of basis for the 4-spinor,

\[\hat{\Psi} = (|S, +\rangle, |P, +\rangle, |S, -\rangle, |P, -\rangle)^T \] (1)

(where \( S \) and \( P \) refer to Na-3s and Bi-6p orbitals, \([\cdots]\)^T indicates transpose and \( \pm \) refer to the \( z \) component of \( J_z \)), the linearized Hamiltonian close to the Dirac node at \( \mathbf{k}^\perp_D \) splits into two 2 × 2 blocks given by

\[H(\mathbf{q}) = v_F \begin{pmatrix} \beta q_z & q_+ & 0 & 0 \\ q_+ & -\beta q_z & 0 & 0 \\ 0 & 0 & -\beta q_z & -q_+ \\ 0 & 0 & -q_+ & -\beta q_z \end{pmatrix}, \] (2)

where \( \beta < 1 \) is the velocity anisotropy and \( \mathbf{q} = \mathbf{k} - \mathbf{k}^\perp_D \).

The upper and lower blocks are compactly written as

\[H_+ = v_F (q_+ q_x - q_z q_y + \beta q_z q_x), \]
\[H_- = v_F (-q_+ q_x - q_z q_y + \beta q_z q_x), \] (3)

where \( \{|\tau\} \) are Pauli matrices acting in the orbital subspace \( (S, P) \). \( H_\tau \) represent Weyl nodes of chirality \( \chi = \mp \) (\( \chi \) is the determinant of the velocity matrix \( \nabla / v_F \) defined by the relation \( H_\tau = \mathbf{q} \cdot \mathbf{V}_\tau / v_F \)).

In zero magnetic field (\( \mathbf{B} = 0 \)), the two Weyl nodes coincide in \( \mathbf{k} \) space. If TRI is broken in finite \( \mathbf{B} \), the
Zeeman energy leads to their separation. A characteristic of Weyl states is that the \( n = 0 \) Landau level is chiral. The separation of each Dirac node into two protected Weyl nodes is shown explicitly in the band calculations in REF.\(^3\). In an electric field \( E \), the chiral anomaly is predicted for carriers occupying the \( n = 0 \) level (FIG. 1).

Symmetry protection of 3D Dirac nodes under the point-group rotation \( C_n \) ultimately derives from the orthogonality of states that transform under different irreducible representations\(^7\). With the choice for the inversion operator \( P = \tau_z \), the distinct eigenfunctions of \( C_n \) protect the Dirac nodes for \( n = 3 \) (NaBi) and \( n = 4 \) (Cd₃As₂). As a consequence, matrix elements formed between states in the conduction and valence bands must vanish. If \( P = \tau_x \) (the case of \( \beta \)-BiO₂), the Dirac node is pinned at a TRIM. The remaining case \( P = I \) has not yet led to a candidate. The process of creating Dirac nodes by enforcing band crossing (say, by tuning hydrostatic pressure \( P \)) has been discussed in detail in REFs.\(^{15,26}\). For a system with broken inversion symmetry, Dirac nodes are stable over a finite interval \((P_r,P_l)\). Starting at \( P_r \), each Dirac node separates into two Weyl nodes, which move apart in \( k \) space to subsequently recombine when \( P_r \) reaches \( P_l \). Again, just as in the case of TRI breaking, the Weyl nodes remain symmetry protected in the interval \((P_r,P_l)\). By contrast, in a system with inversion symmetry, the Dirac node — realized at a single \( P_r \) — is unprotected and unstable. States in the unstable node are not expected to exhibit the properties of Weyl fermions.

**Landau quantization**

**Chiral anomaly.** In a strong magnetic field \( B \| \hat{z} \), the bulk electronic states in a conventional semimetal are quantized into Landau levels (LLs) indexed by \( n = 0, 1, 2, \ldots \). For each of the LLs, the energy disperses as \( E(n,k_z) = (\hbar k_z)^2/2m, \) where \( m \) is the effective mass for dispersion \( ||z \). A distinguishing feature of Weyl nodes is that the LL with \( n = 0 \) (hereafter called LLO) is chiral, such that \( E(0,k_z) = \pm v_Fk_z \) (FIG. 1a). For the Weyl node
with $\chi = +1$, the slope $\partial E/\partial k_x$ is positive, so the velocity $v||B$, whereas for $\chi = -1$, $v||-B$. Hence, when $B$ exceeds $B_0$ (the field at which the chemical potential $\mu$ enters the LL0), we have two independent populations $N_L$ and $N_R$ of massless left-moving and right-moving massless fermions, respectively. Because they do not intermix, their respective current densities $J_L$ and $J_R$ are independently conserved. Hence, the total current density $J = J_L + J_R$ and difference current density $J' = J_L - J_R$ are also conserved, viz.

$$\partial_\rho \rho + V \cdot J = 0, \quad \partial_\rho \rho^5 + V \cdot J^5 = 0,$$

(4)

where $\rho$ and $\rho^5$ are the respective charge densities (superscripts 5 reflect the chirality matrix $\chi$, see the Supplementary Information for details). Equation (4) expresses chiral symmetry of massless fermions in the LL0.

Application of an electric field $E||B||z$ causes the population of the branch moving down the potential incline (say, $N_L$) to increase, while the uphill branch decreases at the same rate — the chiral symmetry is broken by coupling to $E$ and $B$. The pumping rate ($dN^5/dt$) is the product of the 2D density of states $T^5 = L^2/(2\pi e^2)$ of the LL0 and the rate of increase of available states driven by $E$, $dn/dt = (L/2\pi)\delta k_f/d\epsilon$ (where $L^3$ is the sample volume and $\epsilon_f = \hbar/eB$ the magnetic length).

We have

$$\frac{1}{L^3} \frac{dN^5}{dt} = \frac{1}{2\pi e^2} \frac{1}{2\pi} \frac{dE}{\hbar} = \frac{e^2}{4\pi^2 \hbar^2} E \cdot B \equiv A.$$

(5)

The anomaly term $A$ acts like a source term that ruins conservation of the current $J^5$, which we express as

$$\partial_\rho \rho^5 + V \cdot J^5 = A,$$

(6)

where $\rho^5 = N^5 e/L^2$. Equation (6), which represents the chiral anomaly in a semimetal, implies the emergence of a new axial charge current that strongly enhances the conductivity $\sigma_{ax}$ if the axial relaxation time $\tau_a$ greatly exceeds $\tau_p$, the conventional transport lifetime. The enhancement is observable as a large, negative LMR ($B||E$). In the weak-field limit, the chiral anomaly leads to the striking prediction\textsuperscript{15} that the longitudinal conductivity should increase quadratically with $B(||z)$ as

$$\sigma_{ax} = \frac{e^2}{4\pi^2 \hbar^2} \frac{v (eB)^2 v^2}{c \mu^2} \tau_a,$$

(7)

where $v$ is the band velocity, $c$ the velocity of light and $\tau_a$ is the axial lifetime. Fits to this expression have been used to estimate $\tau_a$.

**Chiral magnetic effect and zero sound.** The CME\textsuperscript{21,28,29} and the chiral anomaly are closely related. The CME describes the appearance of an axial charge current in Weyl fermions. This current is either $||B$ or $||-B$ (depending on $\chi$) in a magnetic field $B$. In a Weyl semimetal, the current is given by $J = \frac{\mu}{2}\mu_0 B$, where $\mu = \mu_R - \mu_L$ is the difference of the chemical potentials $\mu_R$ and $\mu_L$, at the two Weyl nodes. At equilibrium, the axial current vanishes\textsuperscript{30}. However, in applied $E$, $\mu_R$ is rendered non-zero. The CME current is then finite and gives rise to the non-conservation of $J^5$, as in Eq. (6)\textsuperscript{15}.

The CME also has experimental consequences when the system is periodically driven off equilibrium at high frequencies. It has been predicted that, in a Weyl semimetal with two or more pairs of Weyl nodes, there exists a collective mode in which each Weyl Fermi surface executes a breathing mode with a specific phasing between nodes\textsuperscript{32}. Unlike a plasmon, this mode displays a gapless, linear dispersion because local charge densities and currents rigorously vanish everywhere at all times. This collective mode is referred to as chiral zero sound. The simplest example is sketched in Fig. 1d. The two pairs of Weyl nodes are symmetrically arrayed in the $k_z$–$k_y$ plane with $B||k_z$. As shown, nodes with positive chirality ($\chi = 1$, blue) disperse with velocity $v||B$, whereas nodes with $\chi = -1$ (pink) disperse with $v||-B$.

At an instant in the breathing cycle, the local chemical potential $\mu_B$, lies above $E_p$ in the two nodes shown with larger $k_y$ in Fig. 1d, whereas $\mu_R$ lies below $E_p$ for the nodes with smaller $k_y$. The occupation factor in each node is indicated by thick lines in the accompanying dispersion sketches. With this phasing, it is clear that the charge currents cancel pairwise between the four nodes. Deviations of the charge density from equilibrium also cancel. The cancellations allow the mode to propagate as an acoustic wave. The chiral zero sound may contribute strongly to both the heat capacity and the thermal conductivity at low $T$.

A non-local transport experiment to detect the chiral anomaly in Dirac–Weyl semimetals has been proposed\textsuperscript{33}. The test assumes that an unbalance in the electrochemical potentials $\mu_{E_L}$ of two Weyl nodes (labelled R and L) is established by injection of a pump current at one end $(x = 0)$ of a long thin film sample in an applied probe field $B_p$ (Fig. 1e). The electrochemical potential difference, $\delta \mu_{E_L}(x) = \mu_{E_L}^R(x) - \mu_{E_L}^L(x)$, decays exponentially along the sample’s length as $\delta \mu_{E_L}(x) = \delta \mu_{E_L}(0) e^{-x/\ell_c}$. The diffusion length $\ell_c$ is given by $\ell_c = c/\sqrt{D \tau_c}$, where $D$ is the carrier diffusion constant and $\tau_c$ is the intervalley scattering lifetime. The electrochemical potential unbalance may be detected by pairs of voltage contacts at various points along the sample. Importantly, the chiral current direction of either Weyl node depends on the local direction of the magnetic field. Hence, if one could apply a detection field $B_d$ at the voltage contacts distinct from $B_p$, one could verify if the detected signal has a sign dictated by $B_d \cdot B_p$.

**Artefacts in longitudinal magnetoresistance**

In high-mobility Weyl semimetals, LMR experiments are greatly complicated by artefacts caused by current jetting\textsuperscript{16,19}. We assume that the current density $J$ (spatially averaged over the sample) is $||B||k$. The cyclotron motion of the carriers in the $y$–$z$ plane causes the transverse conductivities $\sigma_{yy}$ and $\sigma_{yz}$ to decrease as $1/(\mu_B B)^2$, whereas the longitudinal conductivity $\sigma_{xx}$ is unaffected. The anisotropy results in pronounced concentration of $J$ into a narrow jet aligned with $B$ and a corresponding reduction at the edges parallel to $k$. Hence, the voltage
drop detected at the edge decreases, masquerading as a negative LMR, even though $\sigma_{\text{intr}}$ is unchanged.

In the low-mobility semimetals, Na$_3$Bi and GdPtBi (in which $\mu \approx 1,000$–2,600 cm$^2$ (Vs)$^{-1}$), the effects of current jetting are observable, but they introduce a relatively weak distortion that can be corrected for. However, in the high-mobility semimetals, TaAs, NbAs and NbP ($\mu > 100,000$ cm$^2$ (Vs)$^{-1}$), the fractional change in the observed negative LMR signal is typically very small (0.1–2%) and confined to weak $B \approx \pm 0.5$ T. At larger $B$, the magnetoresistance becomes strongly positive. Frequently, changing the voltage contact placements reverses the sign of the low-$B$ LMR, suggestive of highly inhomogeneous current-density distribution, and, hence, an extrinsic origin for the observed LMR.

Using numerical simulations of $J(x)$ in a longitudinal $B$, a 'squeeze' test was devised to reveal when current-jetting artefacts present a serious concern$^{20}$. The test is based on simultaneous measurements of the maximum and minimum values attained by $J(x)$, which verifies that both decrease with increasing $B$. The inferred intrinsic curve $\rho_{\text{intr}}$ increases steeply with decreasing $B$ and with increasing $T$, the thickness of a thin, square plate with side $w \gg t$, the thickness of a thin, square plate with side $w \gg t$.

In the low-mobility semimetals, Na$_3$Bi and GdPtBi, $R_{\text{edge}}$ and $R_{\text{spine}}$ both decrease with increasing $B$, which verifies that the decrease in $\rho_{\text{intr}}$ is intrinsic and current-jetting effects are subdominant$^{20}$.

In contrast, $R_{\text{edge}}$ and $R_{\text{spine}}$ in high-mobility Bi show divergent field profiles; even at 100 K, $R_{\text{edge}}$ decreases to values very close to zero, while $R_{\text{spine}}$ increases steeply by a factor of 100 (see FIG. 2a in REF$^{20}$). At 100 K, the negative LMR signal arising from the chiral anomaly is expected to be strongly suppressed because of thermally excited carriers occupying high-lying LLs (as verified in Na$_3$Bi and GdPtBi). The divergent trends in Bi are artefacts arising from strong focussing of the jet along the spine. The test has also been applied to TaAs (see the

**Fig. 2** | Magnetoresistance of Na$_3$Bi and GdPtBi. a | Longitudinal magnetoresistance (LMR) curves measured in Na$_3$Bi in parallel fields $B || \{110\}$ at selected $T$ from 4.5 to 300 K. Below ~100 K, the steep decrease of the longitudinal resistivity $\rho_{\text{intr}}$ in increasing $B$ is direct evidence for the chiral anomaly. b | Measured field profiles from the squeeze test to distinguish chiral anomaly longitudinal magnetoresistance from current-jetting effects. Field profiles of $R_{\text{intr}}(B)$ in Na$_3$Bi with $B || \{110\}$ measured in applying the squeeze test. If current-jetting artefacts are dominant, concentration of $J$ along the spine should lead to a steep increase in $R_{\text{spine}}$ versus $B$ (as observed in pure Bi). Here, $R_{\text{spine}}$ is observed to decrease instead. c | Plots of $R_{\text{spine}}(B)$ and $R_{\text{intr}}(B)$ versus $B$ measured in Na$_3$Bi at 2 K. The inferred intrinsic curve $R_{\text{intr}}$, sandwiched between them, shows that, in the quantum limit ($B > 6$ T), $R_{\text{intr}}$ is 10 times smaller than its value at zero $B$. d | Curves of $\rho$ versus $B$ at selected $T$ from 6 to 200 K (with $J$ and $B || \{110\}$) of half-Heusler GdPtBi. e | $R_{\text{spine}}$ at $T$ from 2 to 100 K of GdPtBi. f | The profiles of $R_{\text{spine}}$ and $R_{\text{intr}}$ at 2 K and the inferred $R_{\text{intr}}$ of GdPtBi. Panel a reprinted with permission from REF$^{14}$, AAAS. Panels b, c, e and f reprinted from REF$^{20}$, CC BY 4.0 (https://creativecommons.org/licenses/by/4.0). Panel d reprinted from REF$^{20}$, Springer Nature Limited.
section titled ‘Weyl semimetals’). A rule of thumb is to compare \(B_{\text{in}}\) with \(B_{\text{out}}\) (the field at which \(\mu_B E\) exceeds \(\sim 5\)). If \(B_{\text{in}} < B_{\text{out}}\), the steep growth of current-jetting artefacts effectively precludes reliable observation of any intrinsic LMR in the quantum limit.

**Needle-shaped samples.** There exists a proposal\(^{31}\) that recommends using samples with very high aspect ratio (needles) to mitigate current-jetting effects. This proposal merits careful re-examination. Solutions of the Laplace equation in finite \(B\) are scale invariant. Since there is no intrinsic length scale, current jetting should remain visible in a high-mobility sample until the needle diameter is reduced to scales relevant in the quantum transport regime. A characteristic in needles is that, for the artefactual negative LMR to appear, \(B\) must be aligned exquisitely close to the needle axis. The angular width of the artefactual negative LMR may be less than 1°. Hence, seeing such small angular widths raises a red flag for current jetting.

**Planar Hall effect.** The planar Hall effect is a transverse voltage \(V_{\text{xy}}\) observed when an in-plane \(B\) is at an angle \(\theta\) to \(\langle J \rangle\). \(V_{\text{xy}}\) is given by \((\rho_{xx} - \rho_{yy}) \sin \theta \cos \theta\), where \(\rho_{xx}\) and \(\rho_{yy}\) are resistivities measured parallel and perpendicular to \(B\), respectively. As \(V_{\text{xy}}\) is symmetric in \(B\) (non-Onsager), it is an off-diagonal angular magnetoresistance signal rather than a true Hall effect. The effect arises whenever the principal axes of the tensor \(\rho\) are locked to \(B\) instead of the crystal axes. An example is the field-induced anisotropy in the conductivity tensor in any semimetal, as shown in pure Bi\(^{20}\) anisotropy in the conductivity tensor in any semimetal, where \(\chi_1\) is the principal axes of the tensor rather than a true Hall effect. The effect arises whenever the spin–orbit coupling causes the axial lifetime \(\tau\) to be much longer than the Drude lifetime \(\tau_D\) at a low mobility. The planar Hall effect is a transverse voltage \(V_{\text{xy}}\) observed when an in-plane \(B\) is at an angle \(\theta\) to \(\langle J \rangle\). \(V_{\text{xy}}\) is given by \((\rho_{xx} - \rho_{yy}) \sin \theta \cos \theta\), where \(\rho_{xx}\) and \(\rho_{yy}\) are resistivities measured parallel and perpendicular to \(B\), respectively. As \(V_{\text{xy}}\) is symmetric in \(B\) (non-Onsager), it is an off-diagonal angular magnetoresistance signal rather than a true Hall effect. The effect arises whenever the principal axes of the tensor \(\rho\) are locked to \(B\) instead of the crystal axes. An example is the field-induced anisotropy in the conductivity tensor in any semimetal, as shown in pure Bi\(^{20}\) anisotropy in the conductivity tensor in any semimetal, where \(\chi_1\) is the principal axes of the tensor rather than a true Hall effect.

**Magnetoresistance of semimetals**

**The Dirac semimetal \(\text{Na}_3\text{Bi}.** The Dirac semimetal \(\text{Na}_3\text{Bi}\) crystallizes in the hexagonal \(P6_3/mmc\) phase \((D_{3h}^5)\). The conduction band is primarily derived from the Na-3s states, while the uppermost valence band is derived from Bi-6p\(_{xy}\) states\(^{17}\). The strong spin–orbit coupling causes the Na-3s band to lie below Bi-6p\(_{xy}\) by \(\sim 0.7\) eV. In addition, spin–orbit coupling lifts the heavy-hole band \(|P, \pm \frac{3}{2}\rangle\) above the light-hole band \(|P, \pm \frac{1}{2}\rangle\). The resulting band crossings lead to Dirac nodes at \(k^3_\parallel = (0, 0, \pm 0.26\pi/c)\) along the \(z\) axis (Γ–\(A\)). Because the \(S\) and \(P\) bands transform under \(C_4\) with different irreducible representations and TRI is preserved, the nodes are protected against gap formation. At energies near the node energy \(E_N\) it is sufficient to retain the states

\[
\begin{align*}
\left| S^+, \pm \frac{1}{2}\right\rangle, \\
\left| P^+, \pm \frac{3}{2}\right\rangle
\end{align*}
\]

which are bonding and antibonding combinations, respectively, of orbitals centred at the two Na ions (for \(S\)) and Bi ions (for \(P\)) with parity eigenvalues \(\pm\).

Applying a \(B\) field breaks TRI and splits each Dirac node into two Weyl nodes of opposite chirality \(\chi\). Under Landau quantization, the LL0 is strictly chiral, dispersing either \(||B||\) or \(-|B|\), depending on \(\chi\). The existence of only two Dirac nodes with \(E\) very close to \(E_0\) and the absence of other (spectator) bands at \(E_0\) make \(\text{Na}_3\text{Bi}\) a very attractive platform to search for the chiral anomaly, despite its hypersensitivity to moist air. In samples with low Na vacancies, \(E_0\) lies just below \(E_N\). As \(B\) decreases from 300 K, the resistivity \(\rho\) rises monotonically by a factor of \(\sim 20\) to saturate at 21 m\(\Omega\)cm below 20 K, with Hall density \(n_H \sim 1 \times 10^{13} \text{ cm}^{-2}\) and mobility \(\mu_H \sim 2,600 \text{ cm}^2(\text{Vs})^{-1}\)\(^{24}\). The distinctly non-metallic profile is also observed in GdPtBi. The weak Shubnikov–de Haas oscillations observed indicate that the LL0 is entered at \(\sim 6\) T.

When \(B\) is aligned with \(\langle J \rangle||k\), negative LMR becomes apparent at \(\sim 100\) K (Fig. 2a). At 4.5 K, \(\rho_{xy}\) undergoes a sixfold decrease before saturating above 8 T, consistent with the appearance of the chiral anomaly\(^{34}\). The low mobility \((2,600 \text{ cm}^2(\text{Vs})^{-1})\) implies that current-jetting artefacts should only appear above 11 T\(^{35}\). However, the curve of \(\sigma_{xy}(B)\) lies above \(\rho_{xy}(B)\). Moreover, \(\rho_{xy}(B)\) displays a broad minimum near 10 T, followed by a gradual increase at larger \(B\) as shown in (Fig. 2b), which reflects the competition between the intrinsic LMR and current-jetting effects that are seen above 10 T. Both features are striking evidence that, despite the low mobility, current-jetting effects can still produce observable distortions. However, by comparing the measured curves against numerical simulations, it is possible to remove the distortions to extract the intrinsic curve \(\rho_{xy}(B)\), which is sandwiched between \(\rho_{xy}(B)\) and \(\rho_{xy}(B)\) (Fig. 2c). The inferred \(\rho_{xy}(B)\) reveals that the chiral anomaly leads to a tenfold decrease in \(\rho_{xy}\) which implies that the axial lifetime \(\tau_a\) is 10× longer than the Drude lifetime \(\tau_D\) at \(B = 0\). This is currently the most reliable measurement of the ratio \(\tau_a/\tau_D\) by D.C. transport. As seen in (Fig. 2c), \(\rho_{xy}\) settles down to a \(B\)-independent value in the LL0 (\(B > 8\) T). This convolution suggests that current-jetting distortions may contribute strongly to the narrowing of the angular width of \(\Delta\sigma_{xx}\). Extension to the more elaborate angular magnetoresistance experiments in tilted \(B\) has not been attempted.

**The half-Heusler \(\text{GdPtBi}.** The unit cell of the half-Heusler \(\text{GdPtBi}\) is comprised of Pt–Gd tetrahedra arrayed in the zincblende structure. The low-lying states together form under \(\Gamma\)–\(3\) and \(\Gamma\)–\(2\) bands trans-
A magnetic field $\mathbf{B}$ lifts the fourfold degeneracy via the Zeeman energy $\hbar \mathbf{B} \cdot \mathbf{J}$. The larger Zeeman gap in $\frac{\sqrt{3}}{2}, \pm \frac{\sqrt{3}}{2}$ (3× that in $\frac{1}{2}, \pm \frac{1}{2}$) leads to band crossings that define Weyl nodes separated in $\mathbf{k}$ space. When $\mathbf{H}$ is aligned with $\mathbf{J}$, a large, negative LMR is observed.5 The curves of $\rho_{xx}$ versus $\mathbf{B}$ at fixed $T$ measured with $\mathbf{B}||\mathbf{J}$ (FIG. 2d) bear a close resemblance to those in Na$_3$Bi. As $T$ is lowered from 200 K, the LMR begins to display a prominent negative trend at 125 K. The curve at 6 K displays a steep decrease of $\rho_{xx}$ by a factor of 5 between 0 and 14 T. The angular dependence of the curves of $\rho_{xx}(B)$ was mapped out in detail for $\mathbf{B}$ tilted within the $x$–$y$ plane, as well as away from it. The enhanced conductivity appears as a broad plume centred around the axis with $\mathbf{B}||\mathbf{J}$.

The tunability of $E_c$ by doping provides an important test of the chiral anomaly that could not be done in Na$_3$Bi. An investigation of 16 samples with $E_c$ on either side of the Weyl node energy (in zero $\mathbf{B}$) demonstrated that the LMR is strikingly large ($\rho$(97K)/$\rho$(0) = 0.1) for samples with $E_c$ close to the Weyl node energy.35 As $E_c$ is moved away from the node energy (as determined by the weak-field Hall effect), the LMR signal is gradually suppressed. These trends are consistent with the chiral anomaly.

The squeeze test on GdPtBi showed that both $R_{\text{spin}}$ and $R_{\text{orb}}$ decreased monotonically with increasing $\mathbf{B}$ ($R_{\text{spin}}(B)$) is shown in FIG. 2e (REF.39). The test confirmed that the negative LMR is intrinsic, just as in Na$_3$Bi. As shown in FIG. 2f, the inferred intrinsic resistance $R_{\text{int}}(B)$ continues to decrease at the highest applied $\mathbf{B}$ (the LLO is reached at $\sim$25 T). To date, Na$_3$Bi and GdPtBi provide the finest evidence (based on LMR) for the chiral anomaly in semimetals.

**The layered semimetal ZrTe$_5$.** The semimetal ZrTe$_5$ crystallizes in a layered structure with space group $Cmcm$ ($D_{4h}^{17}$). Within each $a$–$c$ layer, prismatic chains of ZrTe$_5$ run parallel to the $a$ axis (the needle axis), with adjacent chains linked by additional Te ions. The quasi-2D layers are stacked along the $b$ axis by van der Waals interactions to form the 3D structure. At $\Gamma$, band inversion of orbitals from Te $p$ states leads to a massive Dirac node above a very small gap. Roughly concurrent with the LMR experiment on Na$_3$Bi (REF.34), negative LMR was also reported in ZrTe$_5$ and attributed to the CME (REF.37) (FIG. 3a).

Fig. 3 | Unusual electronic properties of ZrTe$_5$. a) Observation of negative longitudinal magnetoresistance measured at selected $T$ from 5 to 150 K. b) Temperature-dependent evolution of the band structure from 2 to 255 K, measured by angle-resolved photoemission along the direction $\Gamma$–$X$. c) Tilted $\mathbf{B}$ measurement showing the extreme sensitivity of the negative magnetoresistance (MR) to slight tilting of $\mathbf{B}$ out of the layer. The negative component of the MR vanishes for tilt angles $|\theta| > 1^\circ$. In addition to the negative MR, ZrTe$_5$ displays a striking array of Berry curvature effects observable in its Hall resistivity $\rho_{xy}$. d) Anomalous Hall effect (AHE) in a tilted magnetic field $\mathbf{B}$. Curves of the AHE contribution to the Hall resistivity $\rho_{xy}$ are plotted versus $\mathbf{H}$ at selected tilt angles $\theta$ between $\mathbf{B}$ and the $a$ axis. Panel a reprinted from REF.37, Springer Nature Limited. Panel b reprinted from REF.34, CC BY 4.0 (https://creativecommons.org/licenses/by/4.0). Panels c and d reprinted from REF.36, Springer Nature Limited.
Subsequent experiments on ZrTe$_5$ have led to a bewildering array of transport behaviour, caused by the sensitivity of the unusually small carrier population to Te vacancies and variation of the chemical potential $\mu$ with $T$. In early samples grown by chemical vapour transport with high densities of Te vacancies, the resistivity profile $\rho_s$ versus $T$ displays a large peak at the peak temperature $T_p$ (which varies from 60 to 160 K). Angle-resolved photoemission measurements (fig. 3b) on a sample with $T_p = 135$ K show that $\mu$ shifts considerably with $T$ (ref. 38). At 2 K, $\mu$ lies at the bottom of the conduction band close to $\Gamma$. As $T$ is raised to 255 K, $\mu$ crosses the small gap (50–80 meV) to end up near the top of its valence band (at $T_v$, $\mu$ is mid-gap, in agreement with the profile of $\rho_s$). Evidence for a temperature-driven topological transition from a strong topological insulator (TI) to a weak TI occurring at $T_v$ (138 K) has been obtained from infrared spectroscopy 39.

The Te vacancy density is lower in flux-grown crystals. As $T$ decreases from 300 to 2 K, $\rho_s$ rises monotonically to approach saturation below $\sim$10 K. Angle-resolved photoemission spectroscopy measurements 40 show that, at 17 K, $\mu$ lies close to the top of the valence band ($\mu$ does not enter the band gap in the sample studied). The existence of a moderately large ($\sim$50%) negative LMR was confirmed in measurements with $B||\langle J\rangle||a$ (ref. 41). However, when $B$ was tilted at an angle $\theta$ exceeding 1.5°, the negative LMR vanished, as shown in fig. 3c (with $\theta$ defined in the inset). In addition, the LMR monitored at $\theta = 0$ changed to a positive sign when $B$ was increased above 3 T. Both the high sensitivity to $\theta$ and the sign change are not understood.

Apart from the LMR behaviour, ZrTe$_5$ displays a rich assortment of transport features at 2 K engendered by the Berry curvature $\Omega$. As shown in fig. 3d, the Hall resistivity $\rho_{xy}$ displays a large anomalous Hall effect, which has been mapped out over the entire solid angle of the vector $H$ (ref. 42). A very interesting feature is the emergence of a true, anomalous planar Hall effect that reverses sign with the in-plane $H$. The sign reversal with $H$ (Onsager behaviour) distinguishes it from the planar Hall effect engendered by domain wall anisotropy in magnetic thin films.

Under uniaxial stress applied $||a$, $\rho_{xy}$ in flux-grown crystals initially decreases to attain a minimum at a strain $\varepsilon_{\text{min}} \sim 0.12$% and then rises quadratically at higher strain. This unusual behaviour has been interpreted as evidence for a sign change of the mass term $m$ in the Dirac cone arising from a topological phase transition from a strong TI to a weak TI phase 41. A negative LMR that changes sign above $\sim 2$ T and is highly sensitive to $\theta$ was also observed in the two TI phases.

The Dirac semimetal Cd$_3$As$_2$. The Dirac semimetal Cd$_3$As$_2$ has a body-centred tetragonal structure with $C_{4v}^{16}(4\_4\_d)$ symmetry 43. In the unusually large unit cell (with 80 atoms), 1/4 of the Cd sites are vacancies that can attain long-range crystalline order. Despite the large unit cell, the mobility can reach ultrahigh values (10$^9$ cm$^2$ (Vs)$^{-1}$ at 2 K) 44. Unfortunately, the ultrahigh mobility and the large separation between $E_g$ and the Dirac nodes are major impediments to observing the chiral anomaly by LMR experiments. Nonetheless, the successful growth of high-quality thin films has allowed experiments using other approaches (see below). The possible role of the chiral anomaly in planar Hall/angular magnetoresistance effect experiments (in which $B$ is rotated within the plane containing the current and Hall contacts in thin-slab crystals) is explored in nanoplates of Cd$_3$As$_2$ in ref. 45. Differences between the planar Hall or angular magnetoresistance experiments performed on Dirac semimetals and conventional semimetals (Bi) are critically discussed in ref. 20.

Weyl semimetals. Unlike in the Dirac semimetals, the space group of the Weyl semimetals, TaAs, NbAs, TaP and NbP, lacks inversion symmetry. Hence, each Dirac node is already split into isolated Weyl nodes in zero $B$. TaAs belongs to the space group $I4_1\_md$. There exist 24 Weyl nodes, with eight located on the $k_z = 0$ plane and 16 shifted off the plane. The large number of nodes combined with the very high mobilities (100,000 to 150,000 cm$^2$ (Vs)$^{-1}$) make LMR data difficult to analyse. Initially, observations of a small, negative LMR in weak $H$ were interpreted as the chiral anomaly 44,45. However, the LMR observed in the Weyl semimetals is fragile 46,47, changing sign if the voltage contacts are rearranged.

The application of the squeeze test to TaAs demonstrated 3 that current-jetting distortions onset at a field $B_{\text{crit}} \sim 0.5$ T far below $B_q = 7.04$ T ($B_q$ is determined by the quantum oscillations). As $B$ is increased from zero, $R_{\text{min}}$ increases very steeply, whereas $R_{\text{max}}$ fails to values below the detection limit before $B_q$ is attained — the divergent profiles are closely similar to those in Bi (fig. 5a). Based on these findings, it was concluded that LMR experiments cannot be used to confirm the chiral anomaly in TaAs and NbP. Alternative techniques to detect the anomaly-induced currents are described below.

Complementary experiments

Non-local transport. The chiral anomaly may be detected by non-local transport (fig. 1e), as was done in ref. 48. In this study, a focussed ion beam was used to fabricate multiple pairs of probes on a thin-film Cd$_3$As$_2$ sample grown by chemical vapour deposition to measure the non-local signal. The measured non-local signal had contributions from both ohmic diffusion (the conventional current) and the polarization diffusion arising from the charge pumping. The study found that the valley polarization diffusion length is ~3 times the conventional ohmic diffusion length. Because of the short length scales, however, the crucial test of applying an independent $B_\theta$ at the voltage contacts could not be carried out.

Thermopower. The thermopower ($S_{xx}$) provides a probe of the LL0. Instead of the usual quadratic dispersion along $H||k_x$, we now have a strictly 1D linear dispersion. By the Mott formula, the flat density of states causes strong suppression of $S_{xx}$. As noted, the lowest Landau level (LL0) of Weyl fermions is chiral. This striking feature implies that the density of states $D(E)$ is nominally

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Fig. 4 | Experiments on TaAs and GdPtBi. a | In squeeze-test measurements applied to TaAs, the edge and spine resistances, \( R_{\text{edge}} \) and \( R_{\text{spine}} \), respectively, strongly diverge once \( B \) deviates from zero (as shown by the arrows; both curves are displayed on two scales). This implies that the observed LMR is dominated by current-jetting artefacts because of the very high carrier mobility. The inset shows the index plot of Landau levels derived from the weak oscillations in \( R_{\text{spine}} \), which yields \( B_0 = 7.04 \, \text{T} \). b | Curves of the thermopower \( S_{\text{xx}}(B) \) measured in GdPtBi at selected tilt angles \( \theta \) of \( B \) relative to the \( x-y \) plane at \( T = 6.45 \, \text{K} \). The thermal current density \( J_0 \) is applied along the \([110] \) direction. The pronounced decrease of \( S_{\text{xx}} \) with increasing \( B \) at \( \theta = 0 \) is consistent with the increased dominance of the lowest (chiral) Landau level in which \( S_{\text{xx}} \) is strongly suppressed (see text). Weak Shubnikov–de Haas oscillations are observed for the curve at \( \theta = 0 \). c | Raw experimental trace of the temperature gradient \( d\theta/dB \) in the semimetal TaAs as the field \( B \) is swept from \( -9 \) to \( 9 \, \text{T} \) at constant heater power with \( B || J_0 || c \) and \( T \) fixed at \( 2 \, \text{K} \). Giant quantum oscillations are observed in the longitudinal thermal conductivity \( \kappa_{\text{xx}} \). d | The amplitudes of the quantum oscillations in \( \kappa_{c} \) at low \( T \) in TaAs (red curve) are nearly two orders of magnitude larger than the oscillations in the thermal conductivity \( \kappa_{\text{xx}} \) inferred from the longitudinal conductivity \( \sigma_{\text{xx}} \) using the Wiedemann–Franz law (black curve). Panel a reprinted from REF.\(^{52,53,55}\), CC BY 4.0 (https://creativecommons.org/licenses/by/4.0). Panel b reprinted from REF.\(^{35}\), Springer Nature Limited. Panels c and d reprinted from REF.\(^{55}\), CC BY 4.0 (https://creativecommons.org/licenses/by/4.0).
Chiral zero sound. In a recent measurement of the thermal conductivity $\kappa_{xy}$ versus $B$ in TaAs, magnetic quantum oscillations were observed in $\kappa_{xx}$ with remarkably large amplitude when $J_0$ is aligned with $B$ [REF.34] (FIG. 4c). The oscillations have the same period as (but are antiphased with) the Shubnikov–de Haas oscillations in the conductivity $\sigma_{xx}$. The peak-to-peak amplitude $\Delta \kappa_{xx}$~12 W (Km)$^{-1}$ of the largest oscillation is 3.4 times larger than the zero-$B$ value $\kappa$ (FIG. 4d). The overall values of $\kappa_{xx}$ are 50–100 times larger than predicted by the standard Wiedemann–Franz law. After eliminating several potential causes, the authors identify the oscillations as arising from the propagation of the chiral zero sound mode. Both the magnitude and the phase of the oscillations can be fitted to the chiral zero sound model.

Optical experiments. Low-frequency reflectivity experiments on Weyl semimetals can directly probe the enhanced Drude response associated with the chiral anomaly in an applied $B$ [REF.35]. In addition to being immune to current-jetting artefacts, measurements of the Drude enhancement allow the axial relaxation time $\tau_c$ to be estimated.

In addition, pump–probe measurements at THz frequencies have been used to investigate chiral charge pumping and measure carrier relaxation in the Weyl semimetal TaAs in applied $B$ [REF.36]. An intense pump pulse at frequency 3.4 THz with optical field $E_{\text{pump}} \parallel B$ is employed to generate photoexcited carriers. The low photon energy (14 meV) ensures that only carriers very close to the Weyl nodes are excited. The resulting changes to the reflection coefficient are detected by a weak probe pulse. Varying the delay time between pump and probe enables the relaxation time of the excited carriers to be extracted [FIG. 5a]. In addition to the usual hot-carrier contribution (which relaxes very rapidly), there was a long-lived metastable response that persists beyond 1 ns in the presence of $B$ (FIG. 5a). The metastable response is interpreted as evidence for the axial current [FIG. 5b]. This metastable response is present only when $E_{\text{pump}}$ is aligned parallel to $B$, and vanishes when $E_{\text{pump}}$ is perpendicular to $B$. The metastable axial anomaly signal is found to be linear in $B$.

Finally, magnetoteraehertz spectroscopy has been used to measure the real part of the terahertz conductivity spectrum $\sigma$ versus $f$ (frequency) in epitaxial, thin-film Cd$_3$As$_2$ in an applied field $B$ [REF.37]. Utilizing a fast rotation polarizer and synchronous detection, they have measured the complex transmission matrix at frequencies up to 1.6 THz. With the electric field $E_{\text{pump}} \parallel B$, a large enhancement $\Delta \sigma$ is observed. This enhancement is attributed to the chiral anomaly in Cd$_3$As$_2$. The results fit well to the A.C. conductivity expression given in REF.34 (FIG. 5c).

**Perspective**

In this Review, we surveyed experiments on the chiral anomaly in condensed matter physics. Measurement of the LMR is the most direct way to probe the anomaly in Dirac and Weyl semimetals when the field needed to access the LLO ($B_c$) is considerably lower than the onset field $B_{\text{onset}}$ for cyclotronic motion. This is the situation in Na$_3$Bi and GdPtBi. In high-mobility semimetals with $B_{\text{onset}} \gg B_{\text{Qc}}$, current-jetting artefacts present a serious challenge. The squeeze test described is a first step that allows the axial current to be extracted, unencumbered by artefacts. Follow-up experiments will explore its specific properties. We anticipate a new generation of LMR experiments in which multiple probes in cleverly designed configurations allow the current distribution to be mapped in detail. Isolation of the pure axial current without current-jetting artefacts will allow improved tests on the effect of the anomaly on other transport quantities, such as the thermopower and Hall effect, as well as scattering by magnetic versus non-magnetic impurities. The thermal transport and optical reflectivity experiments are much less sensitive to current-jetting.
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Acknowledgements

We are indebted to Stephen Adler for careful reading of the draft and providing valuable comments. N.P. acknowledges the support of the U. S. Army Research Office (ARO contract W911NF-16-1-0116), the U. S. National Science Foundation (grant DMR 1430561) and the Gordon and Betty Moore Foundation’s EPiQS Initiative through grant GBMF4539.

Author contributions

The authors contributed equally to all aspects of the article.

Competing interests

The authors declare no competing interests.

Peer review information

Nature Reviews Physics thanks Chandra Shekhar, Qiang Li and the other, anonymous, reviewer for their contribution to the peer review of this work.

Publisher’s note

Springer Nature remains neutral with regard to jurisdictional claims in published maps and institutional affiliations.

Supplementary information

The online version of the supplementary material is available at https://doi.org/10.1038/s42254-021-00510-9.

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