Signature of the chiral anomaly in a Dirac semimetal – a current plume steered by a magnetic field

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In this talk, we describe recent experimental progress in detecting the chiral anomaly in the Dirac semimetal Na$_3$Bi in the presence of a magnetic field. The chiral anomaly, which plays a fundamental role in chiral gauge theories, was predicted to be observable in crystals by Nielsen and Ninomiya in 1983 [1]. Theoretical progress in identifying and investigating Dirac and Weyl semimetals has revived strong interest in this issue [2–6]. In the Dirac semimetal, the breaking of time-reversal symmetry by a magnetic field $B$ splits each Dirac node into two chiral Weyl nodes. If an electric field $E$ is applied $\parallel B$, charge is predicted to flow between the Weyl nodes. We report the observation in the Dirac semimetal Na$_3$Bi of a novel, negative and highly anisotropic magnetoresistance (MR). We show that the enhanced conductivity has the form of a narrowly defined plume that can be steered by the applied $B$. The novel MR is acutely sensitive to deviations of $B$ from $E$, a feature incompatible with conventional transport. The locking of the current plume to $B$ appears to be a defining signature of the chiral anomaly.

I. PREAMBLE

The intriguing possibility of observing the chiral (or axial) anomaly as a novel charge current in a crystal has been discussed since 1983 [1]. In the field of topological phases of matter, this question has lately received intense interest in the context of Weyl semimetals [2–7].

II. NODE PROTECTION BY SYMMETRY AND WEYL NODES

The problem of how the energy gap in a narrow-gap semiconductor closes and re-opens as a function of a tuning parameter $\mathcal{P}$ (e.g. pressure, doping or temperature) has received considerable attention in the context of topological phases and 3D Dirac states in the bulk [12–18]. The situation in which the gap closes at a single, critical value $\mathcal{P}_c$ used to be regarded as the norm (Fig. 2, a). This situation – dubbed accidental band crossing – is inherently unstable and not very interesting experimentally. Recent progress has shown that Dirac nodes can be protected by symmetry over an extended range of $\mathcal{P}$ (and hence stable) (Fig. 2b). If time-reversal symmetry (TRS) and inversion symmetry (IS) are both present, the protection extends only to Dirac nodes pinned to high-symmetry points (either $\Gamma$ or the corners of the Brillouin Zone) [14]. Unfortunately, this severely limits the number of realistic candidate materials. Later, it was realized that inclusion of point-group symmetry (PGS) extends the protection to nodes away from the high-symmetry points provided they lie on a symmetry axis [15–18]. This crucial insight, which greatly expands the list of materials, led to the prediction by the IOP group [16, 17] that Na$_3$Bi (with the PGS $C_3$) and Cd$_3$As$_2$ ($C_4$) are Dirac semimetals with protected nodes (Fig. 2c). In the past year, ARPES [19–23], STM [24] and transport experiments [25–26] have confirmed these predictions. Recently, surface FS arcs have been detected by ARPES [29].

When a Dirac semimetal is immersed in a strong magnetic field $B$, one of the 3 protecting symmetries, TRS, is removed. Instead of reopening the gap, each of the Dirac nodes is predicted to split into two Weyl nodes [2–6] (Fig. 3A). The Weyl nodes, which may be regarded as...
as monopole sources and sinks of Berry curvature $\vec{F}$ in $k$ space, come in pairs with opposite chiralities $\chi = \pm 1$. It is helpful to regard $\vec{F}(k)$ as an effective magnetic field that lives in $k$-space. If we surround a Weyl node with a Gauss surface $S$, the Chern flux captured defines the chirality $\chi = (1/2\pi) \oint_S \vec{F}(k) \cdot d\vec{S}(k)$. Thus, breaking TRS produces an intense Berry curvature $\vec{F}$ in the vicinity of each Weyl node. Given that $\vec{F}$ imparts a transverse “anomalous” velocity $v_A = \vec{F} \times eE$ to a wave packet, we should expect a host of novel transport features ($E$ is an applied $E$-field). The most interesting involves the chiral anomaly.

III. CHIRAL ANOMALY

A bit of topological physics fell into quantum field theory (QFT) in the late 1960s [8][11]. The charged pions $\pi^{\pm}$ are remarkably long-lived mesons (lifetimes $\tau \sim 2.3 \times 10^{-8}$ s) because, being the lightest hadron, they can only decay by the weak interaction into muons and neutrinos via the processes $\pi^- \to \mu^- + \bar{\nu}_\mu$ and $\pi^+ \to \mu^+ + \nu_\mu$. Mysteriously, the neutral pion $\pi^0$ decays much more quickly (by a factor of 300 million) even though it is a member of the same isospin triplet. Instead of the slow leptonic channels, $\pi^0$ decays by coupling to the electromagnetic field $F_{\mu\nu}$ in the process $\pi^0 \to 2\gamma$. The relevant diagram (Fig. 3B), called the Adler-Bell-Jackiw anomaly [8][9], is a triangular fermion loop that
links the $\pi^0$ (the axial current) to the 2 photons (vector currents) [10] [11]. A hint of the topological nature of the anomaly is that the one-loop diagram receives no further corrections to all orders of perturbation theory. Subsequent research revealed that the anomaly expresses the breaking of a classical symmetry by quantum fluctuations. In modern QFT, the anomaly plays the fundamental role of killing unviable gauge theories [10] [11]. A proposed chiral gauge theory must be anomaly free. Otherwise it is not renormalizable. Arguably the most important example of the anomaly-free rule is the Glashow-Salam-Weinberg electroweak theory, in which the 4 triangle anomalies linking the lepton and quark doublets with gauge bosons sum exactly to zero within each generation. This fortuitous cancellation has been called “magical” [10].

In 1983, Nielsen and Ninomiya (NN) [1] proposed that the chiral anomaly may be observable in a crystal with massless Dirac states in even spacetime dimension (1+1 or 3+1) (Fig. 3C). When we set the mass $m = 0$ in the Dirac equation, the Dirac 4-spinor $\Psi$ describes two independent populations of chiral particles with spin locked either parallel to their momentum (and described by the Weyl 2-spinor $\psi_R$) or antiparallel ($\psi_L$). The chiral charges are separately conserved. This chiral symmetry is ruined as soon as we couple the particles to $F_{\mu\nu}$ (by applying an $E$ field). Chiral charge is pumped from one branch to the other at a rate given by the triangle anomaly, viz.

$$W = \chi \frac{e^3}{4\pi^2 \hbar^2} E \cdot B.$$  

(1)

At first blush, this would appear to be a restatement of charge transport in an ordinary metal. If we regard the left and right-moving states as opposite limbs of a large Fermi surface (FS), the $E$ field simply biases the occupation of the two limbs leading to the familiar FS shift in $k$-space. However, as discussed below, the implied “locking” of $E$ to $B$ is unusual and new. It actually leads to features that fall beyond the purview of standard transport in metals.

Unsurprisingly, the discussion in 2011 of Weyl nodes in the context of 3D topological phases of matter [2] [7] stimulated intense renewed interest in NN’s proposal, initially among theorists but now extending to experimentalists as well. Theoretically, the states in the two Weyl nodes are quantized into Landau levels in strong $B$. A key feature of the Weyl spectrum is that the $n = 0$ Landau level (LL) is chiral, with velocity $v$ parallel or antiparallel to $B$ (Fig. 5A). We have precisely the situation described by NN, except that the 1D dispersion direction is controlled by $B$.

As mentioned, the breaking of time-reversal symmetry (TRS) by $B$ causes each Dirac node to split into two Weyl nodes [15] [17]. In the absence of an electric field $E$, the two chiral populations are separately conserved (Fig. 5A). However, application of $E$ ruins the conservation and leads to a charge-pumping process between nodes, corresponding to the chiral anomaly sketched in Fig. 3B [13] [5] [7]. As shown in Fig. 5A, the lowest Landau level (LL) of the Weyl states in a crystal is chiral. For $\chi = 1$ the lowest LL is right-moving (velocity $v \parallel B$), whereas for $\chi = -1$ it is left-moving. [18] [5] [18] At steady state, $E$ transfers charge from say the $\chi = -1$ branch to the $\chi = 1$ branch at the rate $[5]$

The rate peaks for $E \parallel B$ and vanishes when $E \perp B$. In the weak-$B$ limit, a Berry curvature approach gives
the chiral conductivity \( \sigma _{\chi} \)
\[
\sigma _{\chi} = \frac{e^2}{4\pi^2 \hbar c} \frac{v (\epsilon B_0)^2}{\epsilon_F^2} \tau_v,
\]
where \( \tau_v \) is the intervalley life-time and \( \epsilon_F \) the Fermi energy.

IV. DIRAC SEMIMETAL Na3Bi

The Dirac semimetal Na3Bi grows as mm-sized, deep-purple, plate-like crystals with the largest face parallel to the a-b plane (\( \hat{c} \) is normal to the planes) (Fig. 4k,b). We annealed the crystals for 10 weeks before opening the growth tube. Details of the growth and characterization are reported in Kushwaha et al. [25]. To avoid oxidation, crystals were contacted using silver epoxy in an Argon glove box, and then immersed in paraffin in a capsule before rapid cooling. In Na3Bi, the Dirac nodes are located at the wave vectors \((0,0, \pm k_D)\) with \( k_D \approx 0.1 \text{ Å}^{-1} \) [19] [20]. Initial experiments in our lab [26] on samples with a large Fermi energy \( \epsilon_F \) (400 mV) showed only a positive MR with the anomalous \( B \)-linear profile reported in Cd3As2 [25].

Recent progress in lowering \( \epsilon_F \) has resulted in samples that display a non-metallic resistivity \( \rho \) vs. \( T \) profile and a low Hall density \( n \approx 1 \times 10^{17} \text{ cm}^{-3} \) (Fig. 3C). We estimate the Fermi wavevector \( k_F = 0.012 \text{ Å}^{-1} (8 \times \) smaller than \( k_D)\). The unusual profiles of \( \rho \) and the Hall coefficient \( R_H \) in Panel C imply the zero-\( B \) energy spectrum shown in Panel B. Below \( \sim 10 \text{ K} \), the conductivity is largely due to electrons in the conduction band with electron mobility \( \mu \approx 2,600 \text{ cm}^2/\text{Vs} \). Because the energy gap is zero, holes in the valence band are copiously excited even at low \( T \). As \( T \) rises above 10 K, the increased hole population leads to a steep decrease in \( \rho \) and an inversion of the sign of \( R_H \) at 62 K. From the maximum in \( R_H \) at 105 K, we estimate that \( \epsilon_F \approx 3k_BT \approx 30 \text{ mV} \). As shown in Fig. 5D, the resistivity \( \rho_{xx} \) in a longitudinal field \((\text{B}||\hat{x}), \text{the current})\ displays a remarkable peak at 4.5 K corresponding to a large negative MR (the resistance measured is \( R_{14.23} \) (I applied to contacts 1 and 4, and voltage measured between contacts 2 and 3; the inset in Fig. 5D shows the contact labels and the x and y-axes). Raising \( T \) above \( \sim 100 \text{ K} \) suppresses the peak. The small density \( n \) implies that \( \epsilon_F \) enters the lowest \((N = 0) \) LL at \( B \approx 4-6 \text{ T} \).

V. A NARROW CURRENT PLUME

The axial current is predicted to be large when \( \text{B is aligned with} \ E \). A valuable test then is the demonstration that, if \( E \) is rotated by 90°, the negative magnetoresistance (MR) pattern rotates accordingly, i.e. the axial current maximum is selected by \( \text{B and} \ E \), rather than being pinned to a crystal axis, even in the weak-\( B \) regime.

To test the anisotropy, we rotate \( \text{B} \) in the \( x-y \) plane while still monitoring the resistance \( R_{14.23} \). Figure 4A shows the curves of the resistivity \( \rho_{xx} \) vs. \( B \) measured at 4.5 K at selected \( \phi \) (the angle between \( \text{B and}\ \hat{x} \)). The MR is positive for \( \phi = 90^{\circ} \) \((\text{B}||\hat{y})), \) displaying the nominal \( B \)-linear form observed in Cd3As2 [25] and Na3Bi [20] with \( \text{B}||c \)). As \( \text{B} \) is rotated towards \( \hat{x} \) (\( \phi \) decreased), the MR curves are pulled towards negative values. At alignment \((\phi = 0)\), the longitudinal MR is very large and fully negative (see SI for the unsymmetrized curves as well as results from a second sample).

We then repeat the experiment \textit{in situ} with \( I \) applied to the contacts (3, 5), so that \( E \) is rotated by 90° (the
FIG. 6: Evidence for axial current in Na$_3$Bi (J4) obtained from transport measurements in an in-plane field $B$. Panel A shows plots of the resistivity $\rho_{xx}$ vs. $B$ at selected field angles $\phi$ to the $x$-axis (inferred from resistance $R_{14,23}$, see inset). For $\phi = 90^\circ$, $\rho_{xx}$ displays a $B$-linear positive MR. However, as $\phi \to 0^\circ$ ($B \parallel \hat{x}$), $\rho_{xx}$ is strongly suppressed. Panel B shows plots of $R_{35,26}$ with $E$ rotated by $90^\circ$ relative to Panel A ($B$ makes angle $\phi'$ relative to $\hat{y}$; see inset). The resistance $R_{35,26}$ changes from a positive MR to negative as $\phi' \to 0^\circ$. In both configurations, the negative MR appears only when $B$ is aligned with $E$. The spheres in the insets show the field orientations in 3D perspective. Adapted from Xiong et al. [27].

We remark on the two striking features in Figs. 6A and 6B. The appearance of a strongly varying in-plane MR is rare in metals. The closest example we are aware of is the “anisotropic MR” (AMR) observed in ferromagnets [30–32]. In an in-plane $B$, one observes $\rho(\phi) \sim P + Q \sin^2 \phi$ caused by anisotropic scattering of carriers with velocity $v||M$ vs. $v \perp M$ ($M$ is the magnetization and $P$ and $Q$ are positive parameters). However, as AMR always leads to a positive MR, it cannot account for the very large negative MR observed when $\phi$ (or $\phi'$) $\to 0$. Moreover, no magnetic transitions have been detected in Na$_3$Bi down to 2 K.

FIG. 7: Magnetoresistance of Na$_3$Bi ($R_{14,23}$) when $B$ lies in the $x$-$z$ plane at selected angles $\theta$, with $E \parallel \hat{x}$ axis ($\theta$ is the angle between $B$ and $\hat{x}$). As in Fig. 6, an enhanced current (large negative MR) is observed when $B$ approaches alignment with $E$. When $B$ is perpendicular to $E$, the MR has a positive $B$-linear profile. The field orientations are shown in 3D perspective in the lower insert. From Jun Xiong (unpublished).

VI. STEERING THE PLUME BY FIELD

The second feature is the tight linkage between the MR pattern and the common direction of $B$ and $E$ even in weak $B$ ($\ll 1/\mu \sim 4$ T). If the in-plane MR is caused by FS anisotropies (between $a$ and $b$ axes), the extrema of the oscillation ought to be pinned to the crystal axes. We should not be able to alter the resistivity tensor by rotating the weak $E$ and $B$ fields (this violates linear response). Viewed from ordinary transport, the observed rotation in weak $B$ is indeed anomalous. However, it agrees with the prediction of the chiral anomaly – the axial current peaks when $E$ aligns with $B$ even for weak $B$. Measured resistance is $R_{35,26}$. Remarkably, the observed MR pattern is also rotated by $90^\circ$, even when $B < 1$ T. Defining the angle of $B$ relative to $\hat{y}$ as $\phi'$, we now find that the MR is fully negative when $\phi' = 0$. The curves in Panels A and B are nominally similar, except $\phi = 0$ and $\phi' = 0$ refer to $\hat{x}$ and $\hat{y}$, respectively.
The conductance $G$ of the conductance $40-60$. The scattering rate relaxing the axial current is

$$\sigma = \frac{1}{\tau_e}.$$ 

We form the ratio $\sigma_B = \sigma_{xx}(B, \phi) / \sigma_{xx}(B, 90^\circ)$ vs. $\phi$ as $B$ is tilted in the $x$-$y$ plane at an angle $\phi$ to $\hat{x}$, with $B$ fixed at values $3-7$ T. Figure 8B shows the same measurements but now with $B$ lying in the $x$-$z$ plane at an angle $\phi$ to $\hat{x}$. In both cases, the low-field curves ($B \leq 2$ T) are reasonably described with $\cos^p \phi$ (or $\cos^p \theta$) with $p = 4$ (not shown). However, for $B > 2$ T, the angular widths narrow significantly. Hence, at large $B$, the axial current is observed as a strongly collimated beam in the direction selected by $B$ and $E$ as $\phi$ or $\theta$ is varied.

To see what happens at larger $B$, we extended measurements of $R_{14.23}$ to $B = 35$ T. We observe a new feature at $H_k \sim 23$ T when $B \parallel \hat{y}$. As $B$ is tilted away from $\hat{y}$ ($\phi \rightarrow 55^\circ$), the feature at $H_k$ becomes better resolved as a kink. The steep increase in $\rho_{xx}$ above $H_k$ suggests an electronic instability at large $B$. However, as we decrease $\phi$ below $45^\circ$, $H_k(\phi)$ moves rapidly to above $35$ T. The negative MR curve at $\phi = 0$ remains unaffected by the instability up to $35$ T (the small rising background is from a weak $B_z$ due to a slight misalignment).

To us, the unusual locking of the negative MR pattern in Figs. 8A and 8B to $E$ and $B$ in weak $B$ constitutes strong evidence for the axial current. The experiment confirms the $B^2$ behavior in Eq. 2 and provides a measurement of the long internode scattering lifetime. However, the width of the collimated beam in the direction of $B$ is much narrower than expected from the theory.

In addition to the results here, several groups have also reported observing the chiral anomaly in other materials (primarily TaAs and ZrTe$_5$). The evidence is by and large restricted to the appearance of negative longitudinal MR often bracketed by large positive MR at lower and higher $B$. A concern is that a negative, longitudinal MR restricted to a narrow field interval is (by itself) a rather slender reed to hang a weighty claim from. Further tests, such as the demonstration of the field-steering effect, would appear to be necessary. Nonetheless, the dramatic increase in experimental activity on a growing list of candidate Weyl semimetals is an encouraging sign for the field. We anticipate exciting experimental developments in the next few years.

VII. ANGULAR WIDTH OF PLUME

A surprise to us is the acute sensitivity of the novel current to misalignment at large $B$. We have examined how the conductivity derived from $R_{14.23}$ decays as $B$ is tilted away from $\hat{x}$ in either the $x$-$y$ or the $x$-$z$ plane. Figure 8A displays the curves of $\Delta \sigma_{xx}(B, \phi) = \sigma_{xx}(B, \phi) - \sigma_{xx}(B, 90^\circ)$ vs. $\phi$ as $B$ is tilted in the $x$-$y$ plane at an angle $\phi$ to $\hat{x}$, with $B$ fixed at values $3-7$ T. Figure 8B shows the same measurements but now with $B$ lying in the $x$-$z$ plane at an angle $\phi$ to $\hat{x}$. In both cases, the low-field curves ($B \leq 2$ T) are reasonably described with $\cos^p \phi$ (or $\cos^p \theta$) with $p = 4$ (not shown). However, for $B > 2$ T, the angular widths narrow significantly. Hence, at large $B$, the axial current is observed as a strongly collimated beam in the direction selected by $B$ and $E$ as $\phi$ or $\theta$ is varied.

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