Experimental Tests of the Chiral Anomaly Magnetoresistance in the Dirac-Weyl Semimetals Na₃Bi and GdPtBi

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In the Dirac-Weyl semimetal, the chiral anomaly appears as an “axial” current arising from charge pumping between the lowest (chiral) Landau levels of the Weyl nodes, when an electric field is applied parallel to a magnetic field \( B \). Evidence for the chiral anomaly was obtained from the longitudinal magnetoresistance (LMR) in Na₃Bi and GdPtBi. However, current-jetting effects (focusing of the current density \( J \)) have raised general concerns about LMR experiments. Here, we implement a litmus test that allows the intrinsic LMR in Na₃Bi and GdPtBi to be sharply distinguished from pure current-jetting effects (in pure Bi). Current jetting enhances \( J \) along the mid-line (spine) of the sample while decreasing it at the edge. We measure the distortion by comparing the local voltage drop at the spine (expressed as the resistance \( R_{\text{spine}} \)) with that at the edge (\( R_{\text{edge}} \)). In Bi, \( R_{\text{spine}} \) sharply increases with \( B \), but \( R_{\text{edge}} \) decreases (jetting effects are dominant). However, in Na₃Bi and GdPtBi, both \( R_{\text{spine}} \) and \( R_{\text{edge}} \) decrease (jetting effects are subdominant). A numerical simulation allows the jetting distortions to be removed entirely. We find that the intrinsic longitudinal resistivity \( \rho_{\text{xx}}(B) \) in Na₃Bi decreases by a factor of 10.9 between \( B = 0 \) and 10 T. A second litmus test is obtained from the parametric plot of the planar angular magnetoresistance. These results considerably strengthen the evidence for the intrinsic nature of the chiral-anomaly-induced LMR. We briefly discuss how the squeeze test may be extended to test ZrTe₅.

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I. INTRODUCTION

In the past two decades, research on the Dirac states in graphene and topological insulators has uncovered many novel properties arising from their linear Dirac dispersion. In these materials, the Dirac states are confined to the two-dimensional (2D) plane. Interest in three-dimensional (3D) Dirac states may be traced to the even earlier prediction of Nielsen and Ninomiya (1983) [1] that the chiral anomaly may be observable in crystals (the space-time dimension 3 + 1D needs to be even). The anomaly, which appears as a current in a longitudinal magnetic field \( B \), arises from the breaking of a fundamental, classical symmetry of massless fermions—the chiral symmetry. Recent progress in topological quantum matter has led to several systems that feature protected 3D Dirac and Weyl states in the bulk [2,3].

A crucial step in the search for 3D Dirac states was the realization that inclusion of point group symmetry [with time-reversal (TR) symmetry and inversion symmetry] allows Dirac nodes to be protected anywhere along symmetry axes, instead of being pinned to TR-invariant momenta on the Brillouin zone surface [4,5]. Relaxation of this constraint led to the discovery of Na₃Bi [6] and Cd₃As₂, in which the two Dirac nodes are protected by \( C_3 \) and \( C_4 \) symmetry, respectively. In the absence of \( B \), each Dirac node is described by a \( 4 \times 4 \) Hamiltonian that can be block diagonalized into two \( 2 \times 2 \) Weyl blocks with opposite chiralities (\( \chi = \pm 1 \)). The absence of mixing between the two Weyl fermions expresses the existence of chiral symmetry. In a strong \( B \), the Weyl states are quantized into Landau levels. As shown in Fig. 1(a), a distinguishing feature is that the lowest Landau level (LLL) in each Weyl node is chiral, with a velocity \( \mathbf{v} \) strictly \( ||\mathbf{B}|| \) (or \( -\mathbf{B} \) as dictated by \( \chi \)) [3].

As a result, electrons occupying the LLL segregate into two massless groups—left and right movers, with populations \( N_L \) and \( N_R \), respectively. Independent conservation
of \(N_L\) and \(N_R\) implies that the chiral charge density \(\rho^S = (N_L - N_R)/V\) is conserved, just like the total charge density \(\rho_{\text{tot}} = (N_L + N_R)/V\) \((V\) is the sample volume). However, application of an electric field \(\mathbf{E}\) breaks the chiral symmetry by inducing mixing between the left- and right-moving branches [Fig. 1(a)] (for a pedagogical discussion, see Ref. [7]). A consequence is that conservation of \(\rho^S\) is violated by a quantity \(\mathcal{A}\) called the anomaly term, viz. \(\nabla \cdot \mathbf{J}^S + \partial_\tau \rho^S = e \mathcal{A}\), where \(\mathbf{J}^S\) is the axial current density. [From the density of states in the LLL and the rate of change \(\partial_k/\partial t\) induced by \(\mathbf{E}\), we obtain \(\mathcal{A} = (e^2/4\pi^2\hbar^2) \mathbf{E} \cdot \mathbf{B}\).] The presence of \(\mathbf{J}^S\) is detected as a large, negative, longitudinal magnetoresistance (LMR). This constitutes the chiral anomaly. (The anomaly first appeared in the theory of \(\pi\)-meson decay [8,9]. See Refs. [10,11].) The conditions for observing the anomaly in Dirac semimetals were discussed, e.g., in Refs. [3,12–15].

In 2015, Xiong et al. reported the observation of a fivefold to sixfold suppression of the LMR in Na\(_3\)Bi, identified with the chiral anomaly [16]. A year later, Hirschberger et al. [17] observed the chiral anomaly, including its thermolectric signature in the half-Heusler GdPtBi. Although the low-lying states in GdPtBi are not Dirac-like in zero \(B\), the application of a Zeeman field splits both conduction and valence bands to produce protected crossings that define Weyl nodes. With \(\mathbf{B}||\mathbf{E}\), a fivefold LMR was observed with a profile very similar to that in Na\(_3\)Bi. In both Na\(_3\)Bi and GdPtBi, the carrier mobility is relatively low \((3000\) and \(2000\) cm\(^2\)/Vs at \(2\) K, respectively).

There have also been several reports of negative LMR observed in the Weyl semimetals TaAs, NbP, and analogs [18–21]. However, the weakness of their LMR signals (50–100 times weaker than in Na\(_3\)Bi) and their fragility with respect to the placement of contacts, together with the high mobilities of the Weyl semimetals \((150\) to \(200\) 000 cm\(^2\)/Vs), have raised concerns about current-jetting artifacts [20,21]. As a consequence, there is considerable confusion and uncertainties about LMR experiments, in general, and the LMR reported in the Weyl semimetals, in particular. The concerns seem to have spread to Na\(_3\)Bi and GdPtBi as well, notwithstanding their much larger LMR signal.

There is a good reason for the uncertainties. Among the resistivity matrix elements, measurements of the longitudinal resistivity \(\rho_{xx}\) \((\mathbf{B}||\hat{z})\) are the most vulnerable to inhomogeneous flow caused by current jetting. Even when the LMR signal in a sample is mostly intrinsic, the chiral anomaly produces an intrinsic conductivity anisotropy \(\eta\), which unavoidably produces inhomogeneous current distributions that distort the observed LMR profile. Given the prominent role of LMR in chiral-anomaly investigations, it is highly desirable to understand these effects at a quantitative level and to develop a procedure that removes the distortions.

A major difference between the large LMR systems Na\(_3\)Bi and GdPtBi, on the one hand, and the Weyl semimetals, on the other, is their carrier densities. Because the density is low in both Na\(_3\)Bi \((1 \times 10^{17}\) cm\(^3\)) and in GdPtBi \((1.5 \times 10^{17}\) cm\(^3\)), the field \(B_\parallel\) required to force the chemical potential \(\zeta\) into the LLL is only 5–6 T. By contrast, \(B_\parallel\) is 7–40 T in the Weyl semimetals. As shown in Fig. 1(a), the physics underlying the anomaly involves the occupation of chiral, massless states. Occupation of the higher LLs \((B < B_\parallel)\) leads to strong suppression of the anomaly [17]. Moreover, as discussed below, LMR measurements involve a competition between the anomaly mechanism (“the quantum effect”) and classical current-jetting effects, which onset at a second field scale \(B_{\text{cyc}}\). The relative magnitudes of these field scales dictate which effect dominates.

Here, we report a series of experiments designed to separate intrinsic from extrinsic effects in LMR experiments. Focusing of the current density \(J(r)\) into a beam strongly reduces its value at the edges of a sample. As shown in Sec. II, the effects of current jetting can be neatly factored into a quantity \(L_y\) (line integral of \(J_y\), which can be measured by local voltage contact pairs. By comparing local voltage drops at the maximum and minimum of the profile of \(J\), we devise a litmus test that sharply distinguishes the two chiral-anomaly semimetals from the case of pure Bi (Sec. III). Adopting a quantitative treatment (Sec. IV), we show how the intrinsic \(\rho_{xx}(B)\) can be derived from the local voltage results. Applying this technique to Na\(_3\)Bi, we obtain the intrinsic profiles of \(\rho_{xx}(B)\) and the anisotropy, with current-jetting distortions removed. The degree of distortion at each \(B\) value becomes plainly evident. The competition between the quantum and...
classical effects is described in Sec. V. To look beyond Na$_3$Bi and GdPtBi, we discuss how the tests can be extended using focused ion beam techniques to test ZrTe$_2$, which grows as a narrow ribbon. The Weyl semimetals, e.g., TaAs, require availability of ultrathin films. The planar angular magnetoresistance (Sec. VI) provides a second litmus test—one that is visually direct when displayed as a parametric plot (Sec. VII). In Sec. VIII, we summarize our results.

II. THE SQUEEZE TEST

Current jetting refers to the focusing of the current density $\mathbf{J}(\mathbf{r})$ into a narrow beam $\| \mathbf{B}$ arising from the field-induced anisotropy $u$ of the conductivity (the drift of carriers transverse to $\mathbf{B}$ is suppressed relative to the longitudinal drift). To maximize the gradient of $\mathbf{J}$ along the $y$ axis, we select platelike samples with $L_w \gg t$, where $w$, $L$, and $t$ are the width, length, and thickness, respectively [Fig. 1(b)]. The $x$ and $y$ axes are aligned with the edges and $\mathbf{B} \| \hat{\mathbf{x}}$. As sketched in Fig. 1(c), the profile of $J_x$ vs $y$ is strongly peaked at the center of the sample and suppressed towards the edges. In the squeeze test, we measure the voltage difference across a pair of contacts (blue dots) along the line joining the current contacts (which we call the spine), as well as that across a pair on the edge (yellow dots). To accentuate current-jetting effects, we keep the current contact diameters $d_c$ small ($d_c \ll w$) and place them on the top face of the sample wherever possible. (The squeeze test cannot be applied to needlelike crystals.)

A. Sample preparation

We provide details on the preparation of Na$_3$Bi, which is by far the most difficult of the three materials to work with. Na$_3$Bi crystallizes to form hexagonal platelets with the broad face normal to (001). The crystals investigated here were grown under the same conditions as the samples used in Xiong et al. [16]; they have carrier densities $1 \times 10^{17}$ cm$^{-3}$ and $B_Q$ in the range 5–6 T. These crystals should be distinguished from an earlier batch [22] that have much higher carrier densities ($3–6 \times 10^{19}$ cm$^{-3}$) for which we estimate $B_Q \sim 100$ T. No evidence for negative LMR was obtained in the highly doped crystals [22].

Because of the high Na content, crystals exposed to ambient air undergo complete oxidation in about 5 s. The stainless growth tubes containing the crystals were opened in an argon glovebox equipped with a stereoscopic microscope, and all sample preparation and mounting were performed within the glovebox. The crystals have the ductility of a soft metal. Using a sharp razor, we cleaved the bulk crystal into platelets $1 \times 1 \text{mm}^2$ on a side with thickness $100 \mu m$. Current and voltage contacts were painted on using silver paint (Dupont 4922N). A major difficulty was achieving low-resistance contacts on the top face (for measuring $R_{spine}$). After much experimentation, we found it expedient to remove a thin layer of oxide by lightly sanding with fine emery paper (within the glovebox). The sample was then placed inside a capsule made of G10 epoxy. After sealing the lid with STYCAST epoxy, the capsule was transferred to the cryostat.

We contrast two cases. In case 1, the anisotropy $u = \sigma_{xx}/\sigma_{yy}$ increases in a longitudinal field $\mathbf{B}$ because the transverse conductivity $\sigma_{yy}$ decreases steeply (as a result of cyclotron motion), while $\sigma_{xx}$ is unchanged in $B$. With $\mathbf{B} \| \hat{\mathbf{x}}$, the two-band model gives the resistivity matrix

$$\tilde{\rho}(B) = \begin{bmatrix}
[\sigma_{x} + \sigma_{y}]^{-1} & 0 \\
0 & [\sigma_{y} + \sigma_{x}]^{-1}
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we suppress the $z$ component for simplicity). The zero-$B$ conductivities of the electron and hole pockets are given by $\sigma_e = ne\mu_e$ and $\sigma_h = pe\mu_h$, with $n$ and $p$ the electron and hole densities, respectively, and $e$ the elemental charge. Note that $\mu_e$ and $\mu_h$ are the mobilities in the electron and hole pockets, respectively, and $\Delta_s = (1 + \mu_e^2 B^2)$ and $\Delta_d = (1 + \mu_h^2 B^2)$. With $\mathbf{B} \| \hat{\mathbf{x}}$, the off-diagonal elements vanish. In case 1, we assume that $\sigma_e$ and $\sigma_h$ remain constant. Hence, the observed resistivity $\rho_{xx}$ is unchanged in $B$. However, the transverse conductivity $\sigma_{yy}$ decreases (as $1/B^2$ in high $B$). The anisotropy arises solely from the suppression of the conduction transverse to $\mathbf{B}$ by the cyclotron motion of both species of carriers.

Case 2 is the chiral anomaly regime in the Dirac semimetal. Charge pumping between Landau levels (LLs) of opposite chirality leads to an axial current that causes $\sigma_{xx}$ to increase with $B$. Simultaneously, the 1D nature of the LL dispersion suppresses the transverse conductivity $\sigma_{yy}$. Hence, the increase in $u$ derives equally from the opposite trends in $\sigma_{xx}$ and $\sigma_{yy}$.

Denoting field-induced changes by $\Delta$, we have

$$\Delta u > 0 \Leftrightarrow \Delta \sigma_{xx} > 0; \Delta \sigma_{yy} < 0.$$  

(2)

$$\Delta \sigma_{xx} > 0; \Delta \sigma_{yy} < 0.$$  

(3)

In the test, the voltage drops $V_{spine}$ and $V_{edge}$ are given by

$$V_{edge}(B) = -\rho_{xx}(B) \int_0^x J_z(x, w/2; B) dx \equiv \rho_{xx} \mathcal{L}_z,$$  

(4)

$$V_{spine}(B) = -\rho_{xx}(B) \int_0^x J_z(x, 0; B) dx \equiv \rho_{xx} \mathcal{L}_s,$$  

(5)

where $\mathcal{L}_z(B)$ and $\mathcal{L}_s(B)$ are the line integrals of $J_z$ along the edge and spine, respectively ($\ell$ is the spacing between voltage contacts). The intrinsic $B$ dependence [expressed in $\rho_{xx}(B)$] has been cleanly separated from the extrinsic $B$ dependence of $\mathcal{L}_z(B)$ and $\mathcal{L}_o(B)$, which arises from current focusing effects.
The area under the curve of $J_x$ vs $y$ is conserved, i.e.,

$$\int_{w/2}^{w/2} J_x(x, y; B) dy = I,$$  \hspace{1cm} (6)

with $I$ the applied current. At $B = 0$, we may take $J$ to be uniform with the magnitude $J_0 = I/(wt)$. The line integral reduces to $L_0 = J_0 \ell$. In finite $B$, focusing of the current beam implies that the current density is maximum along the spine and minimum at the edge, i.e., $J_x(x, w/2; B) > J_x(x, 0; B)$. Moreover, Eq. (6) implies that $J_x(x, 0; B) > J_0 > J_x(x, w/2; B)$. Hence, the line integrals satisfy the inequalities

$$L_x(B) > L_0 > L_e(B).$$  \hspace{1cm} (7)

If both $\sigma_e$ and $\sigma_h$ are $B$ independent, as in case 1, we have from Eqs. (4), (5), and (8),

$$V_{\text{spine}}(B) > V_0 > V_{\text{edge}}(B),$$  \hspace{1cm} (8)

where $V_0$ is the voltage drop across both pairs at $B = 0$. Clearly, $V_{\text{spine}}$ increases monotonically with $B$, while $V_{\text{edge}}$ decreases. Physically, focusing the current density along the spine increases the local $E$ field there. Current conservation then requires $J_x$ to be proportionately suppressed along the edges. Measuring $V_{\text{edge}}$ alone yields a negative LMR that is spurious.

In case 2, however, $\rho_{xx}$ decreases intrinsically with $B$ because of the chiral anomaly, while $L_s$ increases. Competition between the two trends is explicitly seen in the profile of $V_{\text{spine}}$ vs $B$. As shown below, in Na$_3$Bi and GdPnBi, the intrinsic decrease in $\rho_{xx}$ dominates, so both $V_{\text{spine}}$ and $V_{\text{edge}}$ decrease with $B$. We remark that, from Eq. (8), $V_{\text{spine}}$ always lies above $V_{\text{edge}}$. Moreover, when the rate of increase in $L_s$ begins to exceed (in the absolute sense) the rate of decrease in $\rho_{xx}$ at sufficiently large $B$, the curve of $V_{\text{spine}}(B)$ can display a broad minimum above which $V_{\text{spine}}$ increases.

Hence, if both $V_{\text{spine}}$ and $V_{\text{edge}}$ are observed to decrease with increasing $B$, the squeeze test provides positive confirmation that the observed LMR is intrinsic. Their field profiles bracket the intrinsic behavior of $\rho_{xx}$. Conversely, if intrinsic LMR is absent (i.e., $\sigma_{xx}$ is unchanged), $V_{\text{spine}}$ and $V_{\text{edge}}$ display opposite trends (the marginal case when the intrinsic LMR is weak is discussed in Sec. IV).

We remark that the current-jetting effects cannot be eliminated by using very small samples (e.g., using nanolithography). As long as we remain in the classical transport regime, the equations determining the functional form of $J(x, y)$ in strong $B$ are scale invariant. Because intrinsic length scales (e.g., the magnetic length $\ell_B$ or the skin depth $\delta$) are absent in classical dc transport, the same flow pattern is obtained on either mm or micron-length scales.

### III. RESULTS OF SQUEEZE TEST

The results of applying the squeeze test on the three systems are summarized in Fig. 2. In panels (a) and (b), we show the voltage drops $V_{\text{edge}}$ and $V_{\text{spine}}$ measured in pure bismuth (sample B1). The signals are expressed as the effective resistances

$$R_{\text{edge}} = \rho_{xx}(B) \cdot L_e(B)/I,$$  \hspace{1cm} (9)

$$R_{\text{spine}} = \rho_{xx}(B) \cdot L_s(B)/I.$$  \hspace{1cm} (10)

The steep decrease in $R_{\text{edge}}$ [panel (a)] illustrates how an apparent but spurious LMR can easily appear when the mobility is very high (in Bi, $\mu_e$ exceeds $10^6$ cm$^2$/Vs at 4 K). Comparison of $R_{\text{edge}}$ and $R_{\text{spine}}$ measured simultaneously shows that they have opposite trends vs $B$. As $\sigma_e$ and $\sigma_h$ are obviously $B$ independent in Bi, $\rho_{xx}$ is also $B$ independent. By Eqs. (9) and (10), the changes arise solely from $L_e$ and $L_s$. Hence, Figs. 2(a) and 2(b) verify experimentally that $V_{\text{edge}}$ and $V_{\text{spine}}$ display the predicted large variations of opposite signs when current jetting is the sole mechanism present (see simulations in Sec. IV).

Next, we consider Na$_3$Bi. In this sample (N1), $R_{\text{edge}}$ below 20 K decreases by a factor of 50 between $B = 0$ and 10 T [Fig. 2(c)]. This is an order of magnitude larger than observed in Ref. [16]. The increase arises from the enhanced current focusing effect in the present contact placement utilizing small current contacts attached to the broad face of the crystal, as well as a larger $u$. In spite of the enhanced jetting, $R_{\text{spine}}$ shows a pronounced decrease in contrast to the case for Bi. The intrinsic decrease in $\rho_{xx}$ dominates the increase in $L_s$ throughout [see Eq. (10)]. Hence, we conclude that there exists a large intrinsic, negative LMR that forces $R_{\text{spine}}$ to decrease, despite the focusing of $J(r)$ along the ridge. Further evidence for the competing scenario comes from the weak minimum at 10 T in the curves below 40 K in Fig. 2(d). As anticipated above, in large $B$, $\rho_{xx}$ approaches a constant because of the saturation chiral-anomaly term. However, $L_s$ continues to increase because the transverse conductance worsens. Consequently, $R_{\text{spine}}$ goes through a minimum before increasing. This is seen in $R_{\text{spine}}$ but absent in $R_{\text{edge}}$.

A feature that we currently do not understand is the large $V$-shaped cusp in weak $B$. At 100 K, the cusp is prominent in $V_{\text{spine}}$ but absent in $V_{\text{edge}}$.

In Figs. 2(e) and 2(f), we show the field profiles of $R_{\text{edge}}$ and $R_{\text{spine}}$ measured in GdPtBi (sample G1). Again, as in Na$_3$Bi, the anomaly-induced decrease in $\rho_{xx}$ dominates the increase in $L_s$, and $R_{\text{spine}}$ is observed to decrease in increasing $B$. The relative decrease in $R_{\text{edge}}$ is larger than that in $R_{\text{spine}}$. Further, $R_{\text{spine}}$ below 10 K shows the onset of a broad minimum above 10 T [Fig. 2(f)], whereas $R_{\text{edge}}$ continues to fall.
of its microscopic origin. Assuming a constant depends only on the conductivity anisotropy functional form of the inhomogeneous current density To start, we note that, once the boundaries are fixed, the strong current-density inhomogeneity induced by jetting.

These features, in accordance with the discussion above, are amenable to a quantitative analysis that yields the intrinsic field profiles of both $\rho_{xx}$ and $u$ (Sec. IV).

IV. THE INTRINSIC LMR PROFILE

The factorization expressed in Eqs. (9) and (10) allows us to obtain the intrinsic field profile of $\rho_{xx}(B)$ in the face of strong current-density inhomogeneity induced by jetting. To start, we note that, once the boundaries are fixed, the functional form of the inhomogeneous current density $J(\mathbf{r})$ depends only on the conductivity anisotropy $u$ regardless of its microscopic origin. Assuming a constant $\rho_{xx} = \rho_0$ (i.e., case 1), we first calculate, by numerical simulation, the effective resistances $R_{\text{edge}}^0(u)$ and $R_{\text{spine}}^0(u)$ over a broad range of $u$. For simplicity, the simulation is performed for a sample in the 2D limit by solving the anisotropic Laplace equation

$$\left[\sigma_{xx}\partial_x^2 + \sigma_{yy}\partial_y^2\right]\psi(x, y) = 0$$

for the potential $\psi(x, y)$ at selected values of $u$. We used the relaxation method on a triangular mesh with Dirichlet boundary conditions at the current contacts [inset in Fig. 3(a)].

Figure 3(a) displays the calculated curves of $R_{\text{edge}}^0$ and $R_{\text{spine}}^0$ in a 2D sample with the aspect ratio matched to that in the experiment on Na$_3$Bi. As expected, with $\rho_{xx}$ set to a constant, the two curves diverge. This reflects the
FIG. 3. Procedure to obtain the intrinsic LMR curve $R_{\text{intr}}$ vs $B$ and the intrinsic anisotropy $u = \rho_{yy}/\rho_{xx}$ from measurements of $R_{\text{edge}}$ and $R_{\text{spine}}$. Panel (a) plots the calculated resistances $R_{\text{edge}}^0$ and $R_{\text{spine}}^0$ vs $u$ with $\rho_{xx} = \rho_0$ (case 1) [obtained by solving the anisotropic Laplace equation (11)] in a 2D sample (matched to $x$ and $y$ dimensions of the sample). The divergent trends reflect the enhancement of $J(r)$ along the spine and its decrease at the edge. The inset shows the triangulation network generated during the simulation. The stubs on the left and right edges are current contacts. Panel (b) plots the calculated template function $F(u)$ in the semilog scale [Eq. (12)]. Equating the measured ratio $G(B)$ to $F(u)$, and using Eq. (13), we derive the intrinsic field profiles of $R_{\text{intr}}$ and $u$ (see Fig. 4).

FIG. 4. The intrinsic field profiles of $\rho_{xx}(B)$ and anisotropy $u(B)$ derived by applying Eqs. (12)–(14). Panel (a) shows the curves of $R_{\text{spine}}$ (black) and $R_{\text{edge}}$ (red) measured in Na$_3$Bi at 2 K. The inferred intrinsic profile $R_{\text{intr}}$ (blue curve) is sandwiched between the measured curves. In this sample (N1), $R_{\text{intr}}$ decreases by a factor of 10.9 between $B = 0$ and 10 T. We note that $R_{\text{spine}}$ displays a weak minimum near 10 T, which results from the competing trends in $\rho_{xx}(B)$ and $L_s(B)$. Panel (b) displays the intrinsic $u(B)$ in Na$_3$Bi derived from Eq. (13). The corresponding curves for GdPtBi (sample G1 at 2 K) are shown in panels (c) and (d). Again, $R_{\text{intr}}(B)$ in panel (c) is sandwiched between the curves of $R_{\text{edge}}(B)$ and $R_{\text{spine}}(B)$. Panel (d) plots the intrinsic anisotropy $u(B)$. In GdPtBi, the Weyl nodes only appear when $B$ is strong enough to force band crossing. In G1, this occurs at 3.4 T ($u$ remains close to 1 below this field).
simultaneous enhancement of the $E$ field along the spine and its steep decrease at the edge caused by pure current jetting.

From the calculated resistances, we form the ratio

$$\mathcal{F}(u) = R^0_{\text{edge}}/R^0_{\text{spine}} = \mathcal{L}_e/\mathcal{L}_s.$$  \hspace{1cm} (12)

The template curve $\mathcal{F}(u)$, which depends only on $u$, is plotted in semilog scale in Fig. 3(b).

Turning to the values of $R_{\text{edge}}(B)$ and $R_{\text{spine}}(B)$ measured in Na$_3$Bi at the field $B$, we form the ratio $G(B) = R_{\text{edge}}/R_{\text{spine}} = \mathcal{L}_e/\mathcal{L}_s$. Although $G$ is implicitly a function of $u$, it is experimentally determined as a function of $B$ (how $u$ varies with $B$ is not yet known). We remark that $G(B)$ and $\mathcal{F}(u)$ represent the same physical quantity expressed as functions of different variables. To find $u$, we equate $G(B)$ to $\mathcal{F}(u)$ in the template curve. This process leads to the equation

$$u(B) = \mathcal{F}^{-1}(G(B)).$$  \hspace{1cm} (13)

from which we determine $u$ given $B$. Finally, because $R^0_{\text{edge}}$ and $R^0_{\text{spine}}$ are known from the simulation, we obtain the intrinsic profile of $\rho_{xx}$ as a function of $B$ using the relations

$$\rho_{xx}(B) = \frac{R_{\text{edge}}(B)}{R^0_{\text{edge}}(u)} \rho_0 = \frac{R_{\text{spine}}(B)}{R^0_{\text{spine}}(u)} \rho_0.$$  \hspace{1cm} (14)

The redundancy (either resistance may be used) provides a useful check for errors in the analysis.

The results of the analysis are shown in Fig. 4 for both Na$_3$Bi and GdPtBi. In Fig. 4(a), we plot $\rho_{xx}(B)$ in Na$_3$Bi [as the intrinsic curve $R_{\text{intr}}(B)$, in blue]. As expected, $R_{\text{intr}}(B)$ is sandwiched between the measured curves of $R_{\text{edge}}(B)$ and $R_{\text{spine}}(B)$. The field profile of the intrinsic anisotropy $u$ is displayed in Fig. 4(b). It is interesting to note that, in Na$_3$Bi at 2 K, the intrinsic conductivity anisotropy increases to 8 as $B$ is increased to 14 T. This engenders significant distortion of $J(r)$ away from uniform flow. The analysis provides a quantitative measure of how current-jetting effects distort the measurements. At 10 T, $R_{\text{spine}}$ is larger than $R_{\text{intr}}$ by a factor of 2.3, whereas $R_{\text{edge}}$ is 4.0 times smaller than $R_{\text{intr}}$. The plots show explicitly how $R_{\text{spine}}(B)$ still decreases (by a factor of about 5) between $B = 0$ and 10 T, despite the enhancement in $\mathcal{L}_s$ caused by current jetting. Here, we see explicitly that this occurs because the intrinsic LMR is so large (decreasing by a factor of 10.9 between 0 and 10 T) that the current squeezing factor is always subdominant. With the procedure described, this subdominant distortion can be removed entirely. The corresponding profiles of $\rho_{xx}(B)$ and $u(B)$ in GdPtBi are shown in Figs. 4(c) and 4(d), respectively. Unlike the case in Na$_3$Bi, a finite $B$ is required to create the Weyl nodes. This occurs at about 3.4 T in G1. Below this field, the system is isotropic ($u$ close to 1).

### V. QUANTUM VS CLASSICAL EFFECTS

From the experimental viewpoint, it is helpful to view the LMR experiment as a competition between the intrinsic anomaly-induced decrease in $\rho_{xx}$ (a quantum effect) and the distortions engendered by current jetting (classical effect). To observe a large, negative LMR induced by the chiral anomaly, it is imperative to have the chemical potential $\zeta$ enter the LLL. The field at which this occurs, which we call $B_Q$, sets the onset field for this quantum effect. By contrast, the distortions to $J$ caused by current jetting onset at the field $B_{cyc}$, which is set by the inverse mobility $1/\mu$. We write $B_{cyc} = \mathcal{A}/\mu$, where the dimensionless parameter $\mathcal{A} = 5–10$, based on the numerical simulations [Eqs. (12) and (13)].

If $B_Q < B_{cyc}$, the LLL is accessed before classical current distortion appears in increasing $B$. This is the situation in the upper shaded region in the $B_Q$ vs $B_{cyc}$ space in Fig. 5(a). The conditions are favorable for observing the chiral anomaly without worrying about classical current jetting. (To be sure, the chiral anomaly itself leads to a large anisotropy $\sigma_{xx}/\sigma_y$ that can distort $J$. However, this is a quantum effect that follows from the chiral anomaly and can be compensated for.) The measured curves of $R_{\text{spine}}$ and $R_{\text{edge}}$ bracket the intrinsic $\rho_{xx}$, which allows the latter to be obtained, as explained in Eqs. (12)–(14). In both Na$_3$Bi and GdPtBi, $B_Q \sim 5–6$ T, whereas $B_{cyc}$ exceeds 30 T. They fall safely within the shaded area.

![FIG. 5. Competition between the intrinsic chiral-anomaly LMR (with onset field $B_Q$) and the extrinsic classical effect of current jetting (onset field $B_{cyc} = \mathcal{A}/\mu$, $\mathcal{A} = 5–10$). Panel (a): In the shaded upper-half region, the system enters the LLL before the current-jetting effects dominate as $B$ increases ($B_Q < B_{cyc}$). Panel (b) shows a hypothetical case close to the boundary $B_{cyc} = B_Q$. In the profile of $R_{\text{spine}}$, an initial decrease is followed by a steep increase above $B_{cyc}$. For systems in the unshaded region ($B_Q > B_{cyc}$), the LLL is entered long after current-jetting effects become dominant. The classical effect effectively screens the quantum behavior. Bulk crystals of Na$_3$Bi and GdPtBi fall safely in the upper half, whereas TaAs and NbAs fall in the lower half. One way to avoid current jetting in the latter is to lower $B_Q$ by gating ultrathin-film samples.](http://example.com/fig5.jpg)
As $B_Q$ approaches $B_{\text{cyc}}$ [the diagonal boundary in Fig. 5(a)], classical current jetting becomes increasingly problematical. In Fig. 5(b), the schematic curve illustrates the trend of how the quantum behavior can be swamped by the onset of current jetting.

Finally, if $B_Q \gg B_{\text{cyc}}$, classical distortion effects onset long before the LLL is accessed. In this case, $R_{\text{spine}}$ and $R_{\text{edge}}$ display divergent trends vs $B$. Even if an intrinsic LMR exists, we are unable to observe it in the face of the dominant (artificial) change in $V_{xx}$ caused by classical current jetting. In the Weyl semimetals TaAs and NbAs, $B_Q \approx 7$ and 40 T, respectively, whereas $B_{\text{cyc}} \sim 0.4$ T (because of their high mobilities). This makes LMR an unreliable tool for establishing the chiral anomaly in the Weyl semimetals. We illustrate the difficulties with $R_{\text{edge}}$ and $R_{\text{spine}}$ measured in TaAs [Fig. 6(a)] and in NbAs [Fig. 6(b)].

In the initial reports, a weak LMR feature (5\%–10\% overall decrease) was observed in TaAs and identified with the chiral anomaly. Subsequently, several groups found that both the magnitude and sign of the LMR feature are highly sensitive to voltage contact placement. We can, in fact, amplify the negative LMR to nearly 100\%. To apply the squeeze test, we have polished a crystal of TaAs to the form of a thin square plate and mounted contacts in the configuration sketched in Fig. 1(b), with small current contacts (about 80 $\mu$m). As shown in Fig. 6(a), $R_{\text{edge}}$ at 4 K (thick blue curve) displays a steep decrease, falling to a value approaching our limit of resolution at 7 T. Simultaneously, however, $R_{\text{spine}}$ (red) increases rapidly. The two profiles are categorically distinct from those in Na$_3$Bi and GdPtBi [Figs. 2(c)–2(f)] but very similar to the curves for Bi [Figs. 2(a) and 2(b)]. Moreover, the weak SdH oscillations fix the field $B_Q$ needed to access the LLL at 7.04 T [inset in (a)]. Since $1/\mu \sim 0.06$ T, we infer that the classical current-jetting effect onset long before the quantum limit is accessed. Hence, TaAs is deep in the right-bottom corner of the phase diagram in Fig. 5(a). The current-jetting effects appear well before TaAs attains the quantum limit at $B_Q$, and it completely precludes the chiral anomaly from being observed by LMR.

Applying the squeeze test to NbAs next, we display the curves of $R_{\text{edge}}$ (black) and $R_{\text{spine}}$ (red) at 4 K in Fig. 6(b). Here, both $R_{\text{edge}}$ and $R_{\text{spine}}$ increase with $B$, but $R_{\text{spine}}$ increases 100 times faster (in the field interval $0 < B < 8.5$ T, $R_{\text{edge}}$ doubles but $R_{\text{spine}}$ increases by a factor of 280). The vast difference in the rate of change is direct evidence for the squeezing of $J(r)$ along the spine, as depicted in Fig. 1(c). Again, with $B_Q \approx 40$ T, we infer that NbAs falls deep in the right-bottom corner of Fig. 5(a). Classical current jetting dominates the LMR.

It is worth remarking that the squeeze-test results do not invalidate the ARPES evidence, showing that TaAs and NbAs are highly sensitive to voltage contact placement. We can, in fact, amplify the negative LMR to nearly 100\%. To apply the squeeze test, we have polished a crystal of TaAs to the form of a thin square plate and mounted contacts in the configuration sketched in Fig. 1(b), with small current contacts (about 80 $\mu$m). As shown in Fig. 6(a), $R_{\text{edge}}$ at 4 K (thick blue curve) displays a steep decrease, falling to a value approaching our limit of resolution at 7 T. Simultaneously, however, $R_{\text{spine}}$ (red) increases rapidly. The two profiles are categorically distinct from those in Na$_3$Bi and GdPtBi [Figs. 2(c)–2(f)] but very similar to the curves for Bi [Figs. 2(a) and 2(b)]. Moreover, the weak SdH oscillations fix the field $B_Q$ needed to access the LLL at 7.04 T [inset in (a)]. Since $1/\mu \sim 0.06$ T, we infer that the classical current-jetting effect onset long before the quantum limit is accessed. Hence, TaAs is deep in the right-bottom corner of the phase diagram in Fig. 5(a). The current-jetting effects appear well before TaAs attains the quantum limit at $B_Q$, and it completely precludes the chiral anomaly from being observed by LMR.

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NbAs are Weyl semimetals. Rather, they demonstrate that the negative LMR results reported to date in the Weyl semimetals fall deep in the regime where classical current-jetting effects dominate.

Figure 5(a) suggests a way to avoid the screening effect of current jetting for the Weyl semimetals. By growing ultrathin films, one may use gating to lower ζ towards zero in the Weyl nodes. This allows the LLL to be accessed at a much lower B₀. Simultaneously, the increased surface scattering of the carriers will reduce µ (hence increase B_{cyc}). By allowing the quantum effect to onset before the classical effect becomes dominant, both trends shift the “operating point” towards the shaded region B₀ < B_{cyc}. The ability to tune B₀ by gating will enable more tests for mapping out the current density distribution. The squeeze test is actually easier to implement using thin-film samples. Because several groups worldwide are attempting to grow platelike semimetals seem quite encouraging.

Because several groups worldwide are attempting to grow platelike shape is optimal). However, using focused ion beam (FIB) techniques, we may envisage sculpting the appearance of a large Berry curvature in applied B. This implies the simultaneous chiral-anomaly candidate materials that do not readily grow and GdPtBi. The FIB technique can be applied to future mapping out the current density distribution. The squeeze test may then be harnessed to deposit voltage contacts for measuring ∆ρ and θ. The tilt produces potential drops V_{xx} and V_{yx} given by

$$V_{xx}/I = \rho_{bb} + \Delta \rho \cos^2 \theta,$$

$$V_{yx}/I = \Delta \rho \sin \theta \cos \theta,$$

where ρ_{aa} and ρ_{bb} are the resistivities measured along axes a and b, respectively, and ∆ρ = ρ_{aa} − ρ_{bb}.

By convention, the transverse voltage V_{yx} is dubbed the “planar Hall effect” even though it is strictly even in B. As V_{yx} does not satisfy the Onsager relation for a true Hall response, this is a misnomer. [In topological matter, the Berry curvature can generate a true in-plane Hall signal that is odd in B and distinct from V_{yx} in Eq. (16).] To avoid confusion, we call V_{yx} the off-diagonal AMR signal and V_{xx} the longitudinal AMR signal.

Generally, the AMR results are not very informative (the same angular pattern is obtained regardless of the microscopic origin of the anisotropy). However, for our problem, we find that the parametric plot of V_{yx} vs V_{xx} provides a litmus test that distinguishes case 1 from case 2.

In case 1, with θ = 0 (B||k), V_{xx} detects ρ_{aa} L_s^1, its “spurious” decrease as ρ_{aa} decreases arises entirely from L_s. In the orthogonal situation θ = π/2, V_{yx} detects ρ_{bb} L_0 (i.e., beam focusing effects are absent). By juxtaposition, the two measurements reveal how u behaves [see Eq. (2)]. This is best shown by plotting V_{yx} against V_{xx}, with θ as the running parameter at a fixed value of the magnitude B. In weak B, the contours describe small loops circling the zero-B point. As B increases, they expand dramatically away from the zero-B point in the direction of increasing V_{xx}. This lopsided expansion (resembling a shock wave) reflects the sharp increase in the resistivity ρ_{bb} measured orthogonal to B [while ρ_{aa} remains unchanged; see Eq. (1)]. Indeed, from Eq. (1), we have, in the high-B limit,

$$\rho_{bb} \to \frac{(\mu_e \mu_h)^2 B^2}{(\sigma_h + \sigma_e)^2}.$$  

Here, ρ_{bb} increases as B² without saturation. Hence, in case 1, we expect the caliper of the contours (given by ∆ρ) to expand without limit as B².

Case 2 yields a qualitatively different parametric plot. In the chiral anomaly regime, our measurements show that ρ_{aa} (captured by V_{xx} at θ = 0) decreases intrinsically with increasing B, while ρ_{bb} (at θ = π/2) increases by roughly the same fraction. The balanced changes lead to closed contours that expand roughly isotropically from the zero-B point. Moreover, the contour calipers ∆ρ approach saturation at large B.
VII. PARAMETRIC PLOTS

As in Sec. III, we compare the planar angular MR results in the three materials, pure Bi, Na$_3$Bi, and GdPtBi. Figure 7(a) displays the angular profiles $\rho_{xx}$ vs $\theta$ at selected field magnitudes $B$ measured in Bi at 200 K with $\mathbf{J} \parallel \hat{\mathbf{x}}$. As $B$ is tilted away from alignment with $\mathbf{J}$ ($\theta = 0$), $\rho_{xx}$ increases very rapidly at a rate that varies nominally as $B^2$. The overall behavior in $\rho_{xx}$ is a very large increase with $B$ as soon as $|\theta|$ exceeds 10°. However, at $\theta = 0$, a decrease in $\rho_{xx}$ of roughly 50% can be resolved. This is the spurious LMR induced by pure current jetting. The off-diagonal signal $\rho_{yx}$ shows the $\sin \theta \cos \theta$ variation described in Eq. (16) ($\rho_{yx}$ is strictly even in $B$).

The corresponding traces of $\rho_{xx}$ and $\rho_{yx}$ measured in Na$_3$Bi at 2 K are shown in Figs. 7(c) and 7(d), respectively. Although the curves for $\rho_{yx}$ are similar to those in Bi, a qualitatively different behavior in $\rho_{yx}$ becomes apparent. At $\theta = 0$, $\rho_{yx}$ is suppressed by a factor of about 7 (when the chiral anomaly appears). In the transverse direction ($\theta = 90^\circ$), the poor conductance transverse to $\mathbf{B}$ in the LLL raises $\rho_{xx}$ by a factor of about 2.5. In terms of absolute magnitudes, the changes to $\rho_{xx}$ are comparable along the two orthogonal directions, in sharp contrast with the case in Bi. This “balanced” growth leaves a clear imprint in the parametric plots. The off-diagonal signal $\rho_{yx}$ displays the same $\sin \theta \cos \theta$ variation as in Bi.

The plots of $\rho_{xx}$ and $\rho_{yx}$ for GdPtBi in Figs. 7(e) and 7(f) also show a concentric pattern. A complication in GdPtBi is that the nature of the Weyl node creation in the field (by the Zeeman shift of parabolic touching bands) is anisotropic.

![Fig. 7](https://example.com/fig7.png)

**FIG. 7.** Planar AMR measurements in Bi, Na$_3$Bi, and GdPtBi. Panels (a) and (b) show the $\theta$ dependence of the diagonal and off-diagonal AMR signals, expressed as effective resistivities $\rho_{xx}$ and $\rho_{yx}$, respectively, in Bi (sample B2) at 200 K. In panel (a), $\rho_{xx}$ decreases with increasing $B$ when $\theta = 0$ (spurious LMR produced by current jetting). When $|\theta|$ exceeds 20°, $\rho_{xx}$ increases steeply with both $|\theta|$ and $B$. In absolute terms, the increase in $\rho_{xx}$ at $\theta = 90^\circ$ far exceeds its decrease at $\theta = 0$. The off-diagonal MR $\rho_{yx}$ follows the standard $\sin \theta \cos \theta$ variation [panel (b)]. Panel (c) shows curves of $\rho_{xx}$ for Na$_3$Bi (sample N2) at 4 K. The decrease in $\rho_{xx}$ at $\theta = 0$ is roughly comparable with the increase at 90°. A similar balanced pattern is observed in GdPtBi (sample G1) measured at 2 K [panel (e)]. In both case 2 systems, the comparable changes in $\rho_{xx}$ measured at 0° and 90° contrast with the lopsided changes seen in Bi (case 1). The off-diagonal $\rho_{yx}$ in both Na$_3$Bi and GdPtBi [panels (d) and (f)] show the standard $\sin \theta \cos \theta$ profile. The profiles in panel (f) are distorted by anisotropic features associated with the creation of the Weyl pockets. The inset in panel (b) shows the sample geometry with $\mathbf{a} \parallel \mathbf{B}$, both at the tilt angle $\theta$ to $\hat{\mathbf{x}}$. 

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Nonetheless, a balanced growth in that is intrinsic is relatively rare. However, the observation which distorts the variation from the $\sin \theta$ oscillations adds a modulation to the off-diagonal curves, Eqs. (2) and (3). nature of mechanisms that increase the anisotropy [panels (a) and (b)] and case 2 directly reflects the distinct pattern as anticipated above. The contrast between case 1 panel (d)] show concentric orbits that expand in a balanced parametric plots in both Na$_2$Bi and GdPtBi, both $R_{\text{spine}}$ and $R_{\text{edge}}$ decrease. Hence, an inspection of the trend in $R_{\text{spine}}$ allows case 1 (here Bi) to be distinguished from case 2, the two chiral-anomaly semimetals. From the factorization implicit in Eqs. (9) and (10), we relate the experimental ratio $G$ to the curve $F$ obtained by numerical simulation. This enables the intrinsic profiles of both $\rho_{xx}(B)$ and $u$ to be obtained from measurements of $R_{\text{edge}}$ and $R_{\text{spine}}$. After removal of the distortions, $\rho_{xx}$ is seen to decrease by a large factor (10.9) between $B = 0$ and 10 T, while $u$ increases by 8. (For simplicity, the numerical simulation was done in the 2D limit, which we judge is adequate for flakelike samples. Obviously, this can be improved by adopting a fully 3D

VIII. CONCLUDING REMARKS

As mentioned in Sec. I, the existence of a negative LMR that is intrinsic is relatively rare. However, the observation of an artifactual decrease $V_{xx}$ in the LMR geometry is a common experience in high-mobility semimetals. To assist in the task of disentangling the rare intrinsic cases from extrinsic cases (mostly caused by current jetting), we described a test that determines the current-jetting distortions and a procedure for removing them. The squeeze test consists of comparing the effective resistance $R_{\text{spine}}$ measured along the spine (line joining current contacts) with that along an edge $R_{\text{edge}}$. In pure Bi, $R_{\text{spine}}$ increases dramatically with $B$, while $R_{\text{edge}}$ decreases. However, in both Na$_2$Bi and GdPtBi, both $R_{\text{spine}}$ and $R_{\text{edge}}$ decrease. Hence, an inspection of the trend in $R_{\text{spine}}$ allows case 1 (here Bi) to be distinguished from case 2, the two chiral-anomaly semimetals. From the factorization implicit in Eqs. (9) and (10), we relate the experimental ratio $G$ to the curve $F$ obtained by numerical simulation. This enables the intrinsic profiles of both $\rho_{xx}(B)$ and $u$ to be obtained from measurements of $R_{\text{edge}}$ and $R_{\text{spine}}$. After removal of the distortions, $\rho_{xx}$ is seen to decrease by a large factor (10.9) between $B = 0$ and 10 T, while $u$ increases by 8. (For simplicity, the numerical simulation was done in the 2D limit, which we judge is adequate for flakelike samples. Obviously, this can be improved by adopting a fully 3D

FIG. 8. Parametric plots of the planar AMR signals. Panel (a) shows the orbits in pure Bi at 200 K (sample B2) obtained by plotting $\rho_{xx}$ vs $\rho_{yx}$ with $\theta$ as the parameter and with $B$ kept fixed. The orbits start at the left (small $\rho_{xx}$) and end at the right (large $\rho_{xx}$). As $B$ increases, the orbits expand without saturation to the right. Panel (b) shows the same plots measured at 100 K. The increased mobility strongly amplifies the lopsided expansion pattern, which resembles a shock wave. The orbits in the parametric plots measured in Na$_2$Bi [Fig. 8(c) and GdPtBi in (d)] are nominally concentric around the point at $B = 0$. The differences between case 1 and case 2 reflect the very different behaviors of $\rho_{xx}$ at the extreme angles $\theta = 0$ and 90°, as discussed in Fig. 7.
simulation.) The yes or no nature of the test based on inspection of $R_{\text{spine}}$, bolstered by the quantitative analysis that removes the subdominant corrections, adds considerable confidence that the chiral-anomaly LMR profiles in Na$_3$Bi and GdPtBi are intrinsic. Moreover, the subdominant distortion factors arising from current jetting can be effectively removed.

In Sec. V, we described LMR experiments as a competition between the intrinsic quantum effect arising from the chiral anomaly and the classical effects of current jetting. The former describes a phenomenon intrinsic to massless chiral fermions. To see it in full force, the applied $B$ should exceed $B_Q$, the field needed to move $\zeta$ into the LLL. This point seems worth emphasizing because in many reports the claimed anomaly seems to appear in weak fields, $B \ll B_Q$. The experimental concern is that once current jetting appears (at the field scale $B_{\text{cy}}$), it inevitably engenders a dominant, negative LMR profile that is extrinsic in origin. The divergent field profiles of $R_{\text{spine}}$ and $R_{\text{edge}}$ provide a strong warning that the LMR profile is then highly unlikely to be intrinsic.

Looking ahead, we discussed in Sec. V how the classical screening effect from current jetting may be avoided by using ultrathin, gateable films of the Weyl semimetals, which may become available in the near future. For the fourth class of chiral anomaly semimetal ZrTe$_5$ [23,24], we propose using a focused ion beam to sculpt platelike samples that are 10 $\mu$m on a side, and applying micro-lithography to attach contacts for the squeeze test.

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