Superconducting and normal-state anisotropy of the doped topological insulator Sr$_{0.1}$Bi$_2$Se$_3$

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Sr$_x$Bi$_2$Se$_3$ and the related compounds Cu$_x$Bi$_2$Se$_3$ and Nb$_x$Bi$_2$Se$_3$ have attracted considerable interest, as these materials may be realizations of unconventional topological superconductors. Superconductivity with $T_c \sim 3$ K in Sr$_{0.1}$Bi$_2$Se$_3$ arises upon intercalation of Sr into the layered topological insulator Bi$_2$Se$_3$. Here we elucidate the anisotropy of the normal and superconducting state of Sr$_{0.1}$Bi$_2$Se$_3$ with angular dependent magnetotransport and thermodynamic measurements. High resolution x-ray diffraction studies underline the high crystalline quality of the samples. We demonstrate that the normal state electronic and magnetic properties of Sr$_{0.1}$Bi$_2$Se$_3$ are isotropic in the basal plane while we observe a large two-fold in-plane anisotropy of the upper critical field in the superconducting state. Our results support the recently proposed odd-parity nematic state characterized by a nodal gap of Eu symmetry in Sr$_{0.1}$Bi$_2$Se$_3$.

Following the discovery of topological insulators$^1$, the search for a superconducting analogue of a topological insulator has gained considerable interest in the condensed matter physics community. A topological superconductor (TSC)$^3$ has a bulk superconducting energy gap (nodal or nodeless)$^6$ but has gapless surface states which are of great interest both for fundamental physics, as they can host Majorana quasiparticles$^7$, and also for applied physics, as the non-Abelian statistics of surface-state excitations have important implications for robust quantum computing$^8$–$^{10}$.

The topological nature of the superconducting state is determined by the symmetry of the superconducting order parameter and the shape of the Fermi surface. In a time-reversal and inversion symmetric system, odd-parity pairing, where $\Delta(\mathbf{k}) = -\Delta(-\mathbf{k})$, and a Fermi surface that contains an odd number of time-reversal invariant momenta are necessary requirements$^8$. In materials with weak spin–orbit coupling, odd-parity pairing corresponds to spin-triplet pairing; for certain strong spin–orbit coupling systems, unique unconventional superconducting states are possible$^{11}$ that may qualify as topological superconductivity. Currently, two paths towards topological superconductivity are being investigated: proximity-induced TSC$^7$,$^{12}$,$^{13}$ at the interface between a conventional superconductor and a topological insulator or a strong spin–orbit coupled semiconductor, respectively, and via chemical doping of bulk topological insulators. Of the superconducting doped topological insulators, the $M_x$Bi$_2$Se$_3$ family of materials ($M =$ Cu, Nb, Sr)$^{14}$–$^{16}$ has generated the most interest as high quality, mm-scale single crystals are available. Topological order observed via ARPES measurements$^{12}$ and magnetization measurements$^{14,15}$ is consistent with a spin-triplet pairing state. Calorimetry measurements$^{16}$ are not in full agreement with conventional BCS theory, and low-temperature penetration depth measurements$^{17}$ indicate nodes in the superconducting energy gap. The observation of zero-bias conductivity peaks in point-contact spectroscopy measurements$^{20}$–$^{24}$ has been interpreted as evidence for Majorana surface states.

The $M_x$Bi$_2$Se$_3$ family maintains the trigonal $R3m$ structure of the parent compound. Following the surprising observation$^{25}$ of an anomalous Knight shift in Cu$_x$Bi$_2$Se$_3$, demonstrating spin-triplet pairing and twofold anisotropy in the threefold symmetric basal plane, a nematic superconducting state with a two-component order
parameter was proposed. This model accounted for subsequent observations of a spontaneous twofold symmetry below \( T_c \) in several quantities including the upper critical field of Cu\(_x\)Bi\(_2\)Se\(_3\) as determined by calorimetry and of Nb\(_x\)Bi\(_2\)Se\(_3\) as determined by magnetization. This state has \( E_u \) symmetry and odd-parity pairing, and allows for states with complete, albeit anisotropic, superconducting gap as well as for a gap with point nodes. Despite the unconventional nature, the superconducting state has been shown to be robust against disorder scattering.

Figure 1(a) shows the \( R\overline{3}m \) crystal structure of Sr\(_{0.1}\)Bi\(_2\)Se\(_3\), the same as that of the parent compound Bi\(_2\)Se\(_3\) with a slightly extended \( c \) axis due to intercalation of the Sr atom in the gap between adjacent quintuple layers of Bi\(_2\)Se\(_3\), while Fig. 1(b) shows the threefold symmetric basal plane, with the \( a \) (blue) and \( a^* \) (pink) directions marked by arrows. Figure 1(c) shows the proposed twofold symmetric \( \Delta_4 \) superconducting gap structure, which breaks crystallographic rotational symmetry in the basal plane.

Here, we present the first thermodynamic determination of the anisotropy of the upper critical field of Sr\(_{0.1}\)Bi\(_2\)Se\(_3\) single crystals through measurements of the temperature dependence of the reversible magnetization, in addition to angular-dependent magnetotransport measurements. Both quantities yield a large twofold in-plane anisotropy of \( H_{c2} \) in which the high-\( H_{c2} \) direction is aligned with the \( a \) axis of the crystal structure. We find that the normal state resistivity of Sr\(_{0.1}\)Bi\(_2\)Se\(_3\) is isotropic in pairs of samples cut at 90° from the same starting crystal, which excludes conventional mass anisotropy from being the cause of the anisotropy in \( H_{c2} \). Furthermore, temperature dependent measurements of the normal-state magnetization show that Sr\(_{0.1}\)Bi\(_2\)Se\(_3\) is diamagnetic with an isotropic susceptibility of \( \sim -2 \cdot 10^{-6} \) (CGS) which largely originates from the core diamagnetism. X-ray diffraction studies indicate that the extinction rule for the \( R\overline{3}m \) crystal structure is fulfilled to a level of 10\(^{-6}\) implying that deviations from the ideal \( R\overline{3}m \) structure are exceedingly small. We thus conclude that the origin of the twofold anisotropy of the superconducting properties is likely caused by an anisotropic gap structure consistent with the nematic \( E_u \) state.

Results

We present results on a series of Sr\(_{x}\)Bi\(_2\)Se\(_3\) crystals. Bar-shaped crystals #1a and #1b were cut from the same starting piece at 90° with respect to each other. Field-angle dependent resistivity measurements (Fig. 2) reveal that the nematic state is not tied to the current flow direction. Detailed resistivity measurements as function of applied magnetic field, field orientation and temperature (Figs 3 and 4) on crystal #2 yield the anisotropic phase diagram. On crystal #3 we performed magnetic measurements of the nematic state (Fig. 5) and of the normal state susceptibility (Fig. 6). Bar-shaped crystals #4a and #4b, used for resistivity measurements (Fig. 6), were cut from the same starting piece such that they are oriented parallel and perpendicular to the nematic axis, respectively. High-resolution x-ray diffraction on crystal #2 (Fig. 7) reveals high crystalline quality. Finally, we determined the in-plane superconducting anisotropy using magnetization and resistivity measurements on crystal #5, shown in Supplemental Information.
superposition of the two basis gap functions $\Delta_{4}$. Three nematic domains related by rotations of 120° should arise in the sample, giving rise to overall three-field of a superconductor with trigonal symmetry and two-component order parameter has been presented in effective masses $\Gamma$.

Figure 2. $\rho(\theta)$ in an applied magnetic field of $|H_{ab}| = 10$ kG in two crystals of Sr$_{0.1}$Bi$_2$Se$_3$ cut at 90° relative to each other from a single larger crystal in the temperature range 1.7 K (black) to 2.9 K (purple) in increments of 0.1 K. The two crystals were cut from the same starting material at 90° with respect to each other (see inset of Fig. 6). Here, the crystals were intentionally cut such that the long axes did not lie along or perpendicular to the nematic axis. 0° in Fig. 2 marks the crystallographic $a$-axis. The red arrow indicates the direction of current with respect to the crystal axes. The angular dependence of the resistivity thus reflects the angular dependence of the upper critical field $H_{c2}$, as directions with higher $H_{c2}$ will remain superconducting whereas directions with lower $H_{c2}$ will be resistive at a fixed temperature. Twofold anisotropy is evident in both Fig. 2(a,b). As temperature is increased from base temperature through the superconducting transition, the twofold anisotropy is eventually lifted, reaching an angle independent normal state. Our observation that the nematic axis does not depend on the direction of the applied current demonstrates that the observed twofold anisotropy is tied to the crystal structure, and is not an effect due to current flow, such as Lorentz force driven vortex motion.

To further investigate the angular anisotropy of $H_{c2}$, a series of $\rho(T)$ curves were measured on crystal #2 with $T_{c} \approx 2.9$ K in different applied magnetic fields with the field vector along the directions of maximum and minimum in-plane $H_{c2}$ as well as along the $c$-axis of the crystal [Fig. 3(a,b,c)]. X-ray diffraction on this sample (see Fig. 7) reveals that the directions of high (low) in-plane $H_{c2}$ correspond to the crystallographic $a$ and $a'$ directions, respectively (see Fig. 1), consistent with previous reports. Figure 3(a,b,c) show that on increasing field the transitions stay sharp and shift uniformly to lower temperatures. A weak normal-state magnetoresistance is observed only for $H_{c2}/c$. Figure 3(d) shows the magnetic phase diagram along the principal axes with $T_{c}$: the midpoint of the resistive transitions. The in-plane anisotropy $\Gamma$ is $\approx 4.5$. Reported values for the anisotropy (also determined from the resistive midpoints) range from 6.8 for nominal 10% doping to 2.7 for nominal 15% doping, both at 1.9 K, whereas on samples with unspecified doping levels an in-plane anisotropy value of $\approx 2.8$ was obtained in the temperature range covered here, $\Gamma$ is approximately temperature-independent.

The nematic $E_{n}$ state that has been proposed as a possible explanation of the twofold anisotropy of superconducting properties is characterized by a two-component order parameter, which can be expressed as a linear superposition of the two basis gap functions $\Delta_{4}$ and $\Delta_{4}$ (see Fig. 1). An in-depth analysis of the upper critical field of a superconductor with trigonal symmetry and two-component order parameter has been presented in Ref. Three nematic domains related by rotations of 120° should arise in the sample, giving rise to overall threefold symmetry. Instead, the vast majority of reported data including those presented here reveal a simple twofold anisotropy indicative of a single nematic domain. The theoretical analysis reveals that, in contrast to a single component order parameter, a two component order parameter couples linearly to strain fields as parameterized by a coefficient $\delta$, and that such strain fields may serve to pin the nematic vector into a single domain. In particular, for sufficiently strong pinning $\delta$ and near $T_{c2}$, the two component order parameter is effectively reduced to a single component which for $\delta > 0$ is approximately $\Delta_{4}$ and for $\delta < 0$ it is approximately $\Delta_{4}$.

In standard single-band Ginzburg-Landau (GL) theory the anisotropy of $H_{c2}$ is given by the anisotropy of the effective masses:

$$\frac{H_{c2}^{i}}{H_{c2}^{j}} = \frac{m_{i}}{m_{j}} = \Gamma > 1$$

where $m_{i}$ and $m_{j}$ are the effective masses along the principal crystal directions $i$ and $j$. The unit vectors $i, j, k$ define a Cartesian coordinate system; here, $i = a, j = a', k = c$. The angular variation of $H_{c2}$ in the $ij$-plane is then given by:

$$\frac{H_{c2}^{i}}{H_{c2}^{j}} = \frac{m_{i}}{m_{j}} = \Gamma > 1$$

(1)
\[ \theta = \phi + \Gamma - Hc_2 \cos \sin (\theta) \] 

(1)

where \( \theta \) is measured from the \( j \)-direction (low \( Hc_2 \)). Approximating the \( Hc_2 \)-line as linear, the angular dependence of \( Tc(H, \theta) \) in a given field \( H \) can be obtained from Eq. 2 as

\[ Tc(H, \theta) = Tc_0 + \frac{H}{\partial H_c^2(T)/\partial T} \sqrt{\cos^2(\theta) + \Gamma^{-2}\sin^2(\theta)}. \] 

(3)

Data in a field of \( H = 10 \) kG are obtained from the polar diagram of \( \rho(T, \theta) \) [(Fig. 4(a)] by tracing for which values of \( T \) and \( \theta \) the resistivity crosses the 50% value. The results for \( Tc(\theta) \) are shown in Fig. 4(b) together with a fit to Eq. 3. The fit yields an in-plane anisotropy of \( \Gamma \approx 3.8 \), in reasonable agreement with the data in Fig. 3(d). The small difference in anisotropy may arise from deviations from linearity of the phase boundaries.

We obtain the first thermodynamic measurement of the in-plane anisotropy of the upper critical field of \( \text{Sr}_{0.1}\text{Bi}_2\text{Se}_3 \) from the temperature dependence of the magnetization of crystal #3 with \( Tc \approx 3 \) K. Figure 5 shows data taken in a field of 4 kG applied along the high- and low- \( Hc_2 \) directions, respectively. A shift in \( Tc \), defined as the intersection of a linear fit to the \( M(T) \)-data with the \( M = 0 \) line, and a change in the slope \( dM/dT \) with field angle are clearly seen. The inset of Fig. 5 displays the twofold symmetric angular variation of the slope \( dM/dT \) in which a low value of the slope corresponds to a high value of \( Tc \). Such behavior is expected in conventional single-band GL theory of anisotropic superconductors, for which the slope in field direction \( i \) is given as

\[ \frac{\partial M(i)}{\partial T} = -\frac{1}{8\pi\beta_i(\kappa(i)^2)} \frac{\partial H_c^{(i)}}{\partial T} \] 

(4)

where \( H_c^{(i)} = \phi_0/(2\pi\xi_j) \) and \( \kappa(i) = \sqrt{(\lambda_j\lambda_i)}(\xi_j\xi_i) \gg 1 \) are the upper critical field and Ginzburg-Landau parameter in direction \( i \), respectively, and \( \beta_i = 1.16 \) is the Abrikosov number. With the \( T \)-linear variation of \( Hc_2 \) near \( Tc \), one finds that the ratio of the slopes for the high- \( Hc_2 \) and low- \( Hc_2 \) directions is given by the inverse anisotropy
Thus, the data shown in the inset of Fig. 5 indicate an anisotropy of $\Gamma \sim 2$ which is smaller than the value deduced from the resistivity measurements (Fig. 3). Data such as shown in Fig. 5(a) taken over the entire angular range in fields of 4 kG and 6 kG yield the angular dependence of $T_c$ as shown in Fig. 5(b) for 4 kG (blue) and 6 kG (red). Although there is sizable scatter in the data [the error bars in Fig. 5(b) reflect the scatter in $T_c$ obtained on repeated runs], a twofold angular symmetry in this thermodynamic determination of $T_c$ is clearly seen consistent with the twofold symmetry observed in magnetotransport measurements. The data shown in Fig. 5(a) also demonstrate that the superconductivity observed in our Sr$_{0.1}$Bi$_2$Se$_3$ crystals is a bulk phenomenon and not filamentary.

Although transport and magnetization measurements yield similar qualitative features of the superconducting phase diagram of Sr$_{0.1}$Bi$_2$Se$_3$, i.e. a sizable in-plane anisotropy, there are clear quantitative differences in the value of the anisotropy deduced from both techniques. Generally, such differences may arise since magnetization
and resistivity represent different quantities, the expectation value of the magnitude-squared of the superconducting order parameter and the onset of phase coherence across the sample, respectively. Furthermore, the resistively determined phase boundaries depend on the resistivity criterion used; here we use the 50% criterion. Nevertheless, considering that the resistive transitions shown in Fig. 3 appear well-behaved, a difference in anisotropy by a factor of ~2 is surprising. In order to rule out the doping-dependence as a cause of the difference in anisotropy seen in magnetization and magnetoresistance measurements, we performed detailed magnetization and resistivity measurements on a large single crystal, sample #5, shown in Supplemental Information, and reproduce the result that the magnetically determined in-plane superconducting anisotropy is smaller than the resistive result: Γ ~ 2.6 versus 5. The reasons for this unexpected behavior are not understood at present, and may be related to the unusual positive curvature observed in Hc2, in all samples as determined by magnetotransport, or to the existence of surface states which may have different superconducting properties than the bulk.

Figure 6(a) shows the temperature dependence of the resistivity of cross-cut crystals with very sharp superconducting transitions at an onset temperature of ~3.0 K oriented such that the current in crystal #4a flows along the a direction and in crystal #4b along a*, respectively. The anisotropy in the normal-state resistivity is small, <10%. We note that absolute values of the resistivity have an uncertainty of ±15% due to uncertainties in the dimensions of the samples and contact geometry. At the same time, the upper critical field displays a sizable in-plane anisotropy as expressed by the ratio of the effective masses (Eq. 1). For superconductors with essentially isotropic gaps, these effective masses are the same as those entering the normal state conductivity. Our observable sizable in-plane Hc2-anisotropy would imply an in-plane resistivity anisotropy of more than 4, which is clearly not consistent with the data shown in Fig. 6. Furthermore, quantum oscillation measurements on the Nb and Cu homologues suggest that the planar cross-section of the Fermi surface shows little warping, indicating that effective mass anisotropy cannot be the sole cause of the anisotropy in Hc2.

However, for the more general case of anisotropic gaps the GL effective masses are given as

Here, ∫D S denotes an integral over the Fermi surface, N is the electron density, vF is the i-component of the Fermi velocity, and μi is the magnitude of the Fermi velocity, both in general k-dependent. ϕ(k) describes the anisotropy of the gap over the Fermi surface, normalized such that its Fermi surface average is unity. For instance, for a spherical Fermi surface (isotropic normal state electronic structure) and a model gap anisotropy of ϕ(k) = sin(θ) (corresponding to two point nodes on the c-axis) the Hc2-anisotropy for fields applied along the c-axis and for fields applied transverse is 1/2√2. We expect that nodal gap structures with different forms of ϕ(k) and gaps with deep minima will show similar qualitative behavior, namely, that Hc2 measured along the line connecting the nodes (minima) is lower than in a transverse direction. This anisotropy can be expressed in terms of the original two-component model as

where χp, χL, and χVV represent the core diamagnetism, χL+ and χL−, the Pauli paramagnetism and Landau diamagnetism of the conduction electrons, respectively, χCW is the Curie-Weiss paramagnetism, and χVV is the van Vleck diamagnetism and Landau diamagnetism of the conduction electrons, respectively. χCW is the Curie-Weiss contribution, possibly due to magnetic impurities. χp is temperature independent and isotropic, whereas χL+ and χL− are temperature independent but in general anisotropic, depending on the band structure and orbital structure. Since the charge count of Sr0.1Bi2Se3 is much lower than that of typical metals we neglect χp and χL±. With the help of tabulated values, χcore of Sr0.1Bi2Se3 can be estimated as ~2-3.10^-6 (CGS). The observed isotropic diamagnetic response of Sr0.1Bi2Se3 in large part caused by its core diamagnetism, which mainly stems from the Se2− ions. The van Vleck contribution may account for the difference between the measured and expected diamagnetic signals, χVV ~ 0.3 · 10^-6 (CGS). In addition, superimposed onto the diamagnetic signal is a paramagnetic contribution, which approximately follows a Curie-Weiss dependence [Fig. 6(b)]. This contribution is also isotropic, and we attribute it to residual magnetic impurities.

Deviations from the ideal R3m crystal symmetry have been proposed as possible causes of the twofold anisotropy itself or as mechanism of pinning the nematic vector into one domain. We performed x-ray diffraction studies on the crystals used here in order to search for these effects. These measurements revealed a high-degree of structural coherence and phase purity. The determined room-temperature lattice parameters are a = 4.146 Å and c = 28.664 Å, consistent with a rhombohedral R3m crystal symmetry derived from the Sr-intercalated Bi2Sr3 structure. Figure 7 shows I scans centered on various (h, 0, l) zones performed on the same crystal whose transport measurements are shown in Figs 3 and 4. Multiple h values are shown; h = 0 (green circles), h = 2 (blue squares), and h = 2 (pink triangles). At all (h, 0, l) zones examined, only Bragg peaks for which 2h + k + l = 3n is satisfied are observed. This is the extinction rule for the R3m structure. The data shown in Fig. 7 reveal that this extrinc-
The trigonal structure enforces an extinction rule unless $2h + k + l = 3n$, where $n$ is an integer. The allowed peaks show approximately 5 orders of magnitude more intensity than at $l$ values that are not allowed, showing the high quality of the crystal. Any distortions away from a perfect crystal structure would appear as violations of the extinction rule; none are seen. An additional crystal examined shows similar results.

Figure 7. (a) Scans centered at $(h, 0, l)$ for multiple values of $l$ on Sr$_{0.1}$Bi$_2$Se$_3$ crystal #2 used for transport measurements. Multiple values of $h$ are shown; $h = 0$ (green circles), $h = 1$ (blue squares), $h = 2$ (pink triangles). The trigonal structure enforces an extinction rule unless $2h + k + l = 3n$, where $n$ is an integer. The allowed peaks show approximately 5 orders of magnitude more intensity than at $l$ values that are not allowed, showing the high quality of the crystal. Any distortions away from a perfect crystal structure would appear as violations of the extinction rule; none are seen. An additional crystal examined shows similar results. (b) Rocking curve centered at $(h, k, l) = (0, 0, 21)$, showing three closely aligned major grains with a narrow mosaic spread of $\sim 0.04^\circ$.
the same starting crystal demonstrate that the in-plane anisotropy of $H_{c2}$ is tied to the crystal structure and is not induced by the current flow, consistent with the thermodynamic observations. These measurements also show that the normal state resistivity of $\text{Sr}_{0.1}\text{Bi}_2\text{Se}_3$ is isotropic in the plane, thereby excluding conventional effective mass anisotropy as a cause of the $H_{c2}$-anisotropy. Furthermore, temperature dependent measurements of the normal-state magnetization reveal that $\text{Sr}_{0.1}\text{Bi}_2\text{Se}_3$ is diamagnetic with an isotropic susceptibility of $\chi = -2 \times 10^{-7}$ (CGS) which largely originates from the core diamagnetism. These results rule out a possible magnetic origin of the superconducting anisotropy. In addition, x-ray diffraction studies reveal a high degree of structural coherence and phase purity without any detectable deviations from the $\Gamma\delta$ cubic crystal structure that could cause the twofold anisotropy. We thus conclude that the origin of the twofold anisotropy of the superconducting properties is likely caused by an anisotropic gap structure as realized in the nematic $E_\text{state}$. In fact, by specializing the general form of the GL free energy applicable to the two-component $E_z$ order parameter to the $\Delta_{\text{a}}$ and $\Delta_{\text{b}}$ basis functions, we retrieve an anisotropic single-component GL expression that can account for the experimental observations, and indicates the $\Delta_{\text{a}}$ state is selected.

**Methods**

Large high quality single crystals of $\text{Sr}_{0.1}\text{Bi}_2\text{Se}_3$ were grown by the melt-growth technique described in Ref. 38. All crystals regularly showed high volume fraction of superconductivity via magnetic susceptibility measurements with small variation in $T_c$ ranging from 2.9 K to 3.05 K. Thin crystals were cut from as-grown bulk crystals. The material cleaves easily in the basal plane yielding naturally flat surfaces parallel to $aa^*$ in the lattice. Gold contact pads were evaporated, and gold wires were then attached to the crystals using silver epoxy in a conventional 4-point measurement configuration. Typically we cut several samples from the same starting piece as shown for example in the inset of Fig. 6(a). After measuring the first sample the overall orientation of the starting piece is known and we can cut all subsequent samples with approximately known orientation such as mutually perpendicular or at arbitrary angles. Some pairs were aligned parallel and perpendicular to the high-$H_{c2}$ direction [such as in Fig. 6(a)] whereas others were intentionally misaligned (such as in Fig. 2). The crystals were mounted with their long axes parallel to each other such that the angle between current and applied in-plane magnetic field were always the same for both. An AMI 10K$G$ superconducting 3-axis vector magnet was used to apply magnetic field in arbitrary directions without having to physically rotate the sample, and currents smaller than or equal to 1 mA were used for the measurements. Slow rotation of the field direction in $2\pi$ increments ensured thermal equilibrium was maintained. The field was swept clockwise from $0^\circ$ to $400^\circ$ to eliminate any magnetic hysteresis effects. By recording the magnetoresistance during theta scans (a field with fixed magnitude turning from the z-axis of the magnet into the horizontal plane and beyond) for two azimuthal angles, the orientation of the $a-a^*$ plane of the sample with respect to the magnet can be determined with a precision of less than a degree. Such a small level of planar misalignment cannot account for the large observed in-plane anisotropy of $H_{c2}$.

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G.D.G., J.A.S., R.D.Z. sample synthesis. K.W. sample fabrication, magnetotransport measurements, and x-ray diffraction. H.C. magnetization measurements. U.W. magnetization measurements and analysis. A.E.K. and K.W.S., theoretical analysis. Z.I. x-ray diffraction. W.-K.K. experiment design. M.P.S. magnetotransport measurements, analysis, manuscript writing with contributions of U.W. and K.W. All authors reviewed the manuscript.
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