

Large topological Hall effect in a chiral antiferromagnet in hopping transport regime

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The combination of structural chirality and magnetism leads to the formation of spin chirality through noncoplanar magnetic structures, resulting in unusual electronic transport properties. The spin chirality generates nonzero Berry curvature in real space, acting as an emergent magnetic field and contributing to the unconventional anomalous Hall effect, known as the geometrical or topological Hall effect (THE). This study unveils the remarkable occurrence of THE in a chiral antiferromagnetic (AFM) semiconductor EuIr_2P_2 in the hopping regime. It exhibits a complex incommensurately spiral AFM ground state due to its chiral crystalline structure, providing fertile ground for the emergence of topologically nontrivial spin textures such as skyrmions. A substantial THE is observed under finite magnetic fields, making EuIr_2P_2 an exceptional case within the ultralow-conductivity hopping regime for investigating the interplay between topologically nontrivial magnetic structures and hopping carriers. Owing to its semiconducting nature, we have formulated a theoretical model based on Mott's variable range-hopping mechanism, effectively elucidating the temperature and magnetic field-dependent behavior of THE. EuIr_2P_2 thus serves as an ideal candidate for comprehending transport properties in the hopping regime and offers a unique opportunity for the implementation of AFM semiconductor-based spintronic devices.

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I. INTRODUCTION

Chiral crystals exhibit natural absence of inversion centers and mirror symmetry, which is recognized as supporting a wide range of novel physical properties. When magnetism participates into this play ground, noncollinear or noncoplanar magnetic structure is promised by the unique spatial symmetry, enriching the novel physics of the electrons triggered by magnetoelectric field, such as geometrical or topological Hall effect (THE) [1–11]. The research of THE is of fundamental

interest in the discovery of nontrivial magnetic topological state such as skyrmions [1,2,5–7]. Phenomenologically, the THE contributes to the Hall resistivity by manifesting a hump-like feature which is nonproportional to the net magnetization under finite magnetic field and vanishes under high magnetic field. A scalar spin texture, described by $\chi_{ijk} = S_i \cdot (S_j \times S_k)$, where $S_{i,j,k}$ are neighboring spins, generates a real-space Berry curvature (BC) that can be regarded as an emergent magnetic field, resulting in the THE [4]. On the other hand, the THE is also claimed in some magnetic topological materials [12–20]. It is believed that the reciprocal-space BC rising from topological band structures also contributes to the THE. These materials typically fall within the high-conductivity regime ($\sigma > 10^3 \Omega^{-1} \text{cm}^{-1}$), which complicates the speculation of the dominant mechanism of real-space or reciprocal-space BC. To date, research on the THE within the ultralow-conductivity hopping regime ($\sigma < 1 \Omega^{-1} \text{cm}^{-1}$), encompassing both experimental and theoretical investigations, remains notably sparse. The scarcity of materials meeting the criteria of low conductivity, as well as the efficient growth method of high-quality single crystals in large sizes, greatly

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constrains research in the hopping regime. Therefore, it is of fundamental significance to conduct comprehensive research that integrates experimental and theoretical approaches to fill this blank.

Recently, our attention has been drawn to an antiferromagnetic (AFM) semiconductor, EuIr_2P_2 , which might be an excellent candidate with promising properties: it crystalizes in a chiral structure with space group of $P3_221$ [21]; it has an expected energy gap around 400 meV [22]; and it undergoes two AFM phases at $T_{N1} = 5.02(5)$ K and $T_{N2} = 2.94(5)$ K [23]. The previous theoretical calculation claimed a non-collinear magnetic ground state ($T < T_{N2}$) and collinear AFM state ($T_{N2} < T < T_{N1}$) by using a Monte Carlo method in EuIr_2P_2 [22]. However, this simplified model only involved a ferromagnetic (FM) nearest-neighbor interaction and an AFM next-nearest-neighbor interaction of Eu^{2+} ions and neglects the Dzyaloshinskii-Moriya (DM) interaction, which is commonly considered in chiral structures. As the experimental magnetic structure of EuIr_2P_2 is absent, the noncollinear or noncoplanar magnetic structures formed by Eu^{2+} are still expected owing to the chiral structure supported by DM interaction. In this paper, we systematically investigated this compound by magnetoelectric transport, neutron diffraction, and theoretical modeling on high-quality single crystals. Our results demonstrate a nontrivial magnetic structure-induced EuIr_2P_2 , which is rarely seen in the carrier-hopping AFM semiconductors.

II. EXPERIMENTAL DETAILS

A. Sample synthesis and characterization

A well-designed flux method for growing large single crystals of EuIr_2P_2 was performed. Previous reports on this compound have focused on polycrystalline samples obtained by the solid reaction method and tiny single crystals cracked by the arc-melting method, which severely limits the investigation [23]. Here we present a growth method using indium element as flux as well as an oscillatory temperature-control method to promote the growth of single crystals. Starting elements of Eu, Ir, P, and In were put into an alumina crucible in a molar ratio of 1 : 2 : 2 : 30 and then sealed in an evacuated quartz tube. After that the tube was slowly heated to 1373 K, maintained for 10 h, and cooled down to 973 K in an oscillating procedure, as shown in Fig. S1(a) in Supplemental Material (SM) [24]. Hexagonal prism-shaped single crystals with silvery surfaces were yielded after centrifuge. The size of single crystal can be increased to $1 \times 1 \times 3 \text{ mm}^3$ as displayed in the inset of Fig. S1(a) in SM. Energy-dispersive x-ray spectroscopy gives a molar ratio of Eu : Ir : P = 19.61 : 40.73 : 39.66, which is approximately equal to the stoichiometric ratio of this compound.

B. Single-crystal x ray, Laue, and neutron diffraction

The single-crystal x-ray diffraction (SCXRD) was performed by using a self-developed device diffractometer equipped with Mo K_α radiation ($\lambda = 0.71073 \text{ \AA}$). The SCXRD data were then refined by a full-matrix least-squares method on F^2 using the SHELXL-2018/3 program [25]. The crystalline structure refined from SCXRD is summarized in

Tables S1 and S2 in SM. Observed sharp patterns from Laue backward scattering were consistent with the simulated ones, revealing good quality for the single crystals [see Figs. S1(c) and S1(d) in SM]. The single-crystal neutron-diffraction (SCND) measurements were performed at the thermal-neutron diffractometer D23 (ILL, France) on a high-quality EuIr_2P_2 single crystal with mass of 21.6 mg. The neutron wavelength of $1.283 \text{ \AA}$ was used. The momentum transfer $\mathbf{Q} = (H, K, L)$ is defined in hexagonal reciprocal lattice units (r.l.u.) based on the lattice constants $a \approx 6.68 \text{ \AA}$, $c \approx 7.07 \text{ \AA}$, $\alpha = \beta = 90^\circ$, and $\gamma = 120^\circ$, where H, K , and L are Miller indices. The crystal was mounted with the $(H, K, 0)$ reciprocal plane aligned in the horizontal scattering plane. The lifting-counter geometry available at D23 also allows access to many reflections that are out of the $(H, K, 0)$ plane. A standard ILL Orange cryostat was used to cool the sample down to 1.8 K. All SCND measurements were conducted under zero magnetic field.

C. Magnetic and magnetoelectric transport properties

Magnetic properties were measured in a magnetic properties measurement system (MPMS, Quantum Design) via vibrating sample magnetometer (VSM) option. Temperature-dependent susceptibility (2 ~ 300 K) under various magnetic field were measured in zero-field-cooled (ZFC) and field-cooled (FC) configuration. The isothermal magnetization ($-7 \sim 7 \text{ T}$) at various temperatures in four quadrants. Temperature-dependent resistivity, transverse- and Hall resistivity were measured in a physical properties measurement system (PPMS, Quantum Design) via electronic transport option. Data were collected in temperature range of 2 ~ 300 K and magnetic field range from $-7 \sim 7 \text{ T}$ in sweeping mode. Samples were cut and polished with exposed ab plane as well, and a four-terminal configuration of contacts was made on the sample by using the platinum wires which were connected by silver epoxy.

III. RESULTS AND DISCUSSION

Typical side view and top view of the crystalline structure for EuIr_2P_2 is schematized in Fig. 1(a). The helical chains along the c axis formed by Eu atoms are bonded to explicitly illustrate the chiral structure. Figure 1(b) shows magnetic susceptibility of magnetic field parallel to the c axis ($\mu_0 H || c$) and ab plane ($\mu_0 H || ab$) with external magnetic field of $\mu_0 H = 0.01 \text{ T}$. Two magnetic phase transition features were observed at $T_{N1} \sim 5.2 \text{ K}$ and $T_{N2} \sim 3.1 \text{ K}$, which are close to the reported values [23]. For both zero-field-cooling (ZFC) and field-cooling (FC) modes, the measured curves strictly overlapped with each other in the entire temperature range. Thus, only the ZFC data are presented here. Data of both directions under magnetic field of $\mu_0 H = 1 \text{ T}$ in a temperature range from 50 to 300 K could be well fitted via the Curie-Weiss law (see Fig. S2 in SM) [24]. Consequently, the effective magnetic moments were obtained as $7.927(11) \mu_B$ for $\mu_0 H || c$ and $7.970(7) \mu_B$ for $\mu_0 H || ab$, which are consistent with the theoretical value of $7.94 \mu_B$ for Eu^{2+} ($L = 0, J = S = 7/2$). Saturated magnetization [see the inset of Fig. 1(b)] at 2 K for $\mu_0 H || c$ and $\mu_0 H || ab$ attained values close to $7 \mu_B$,

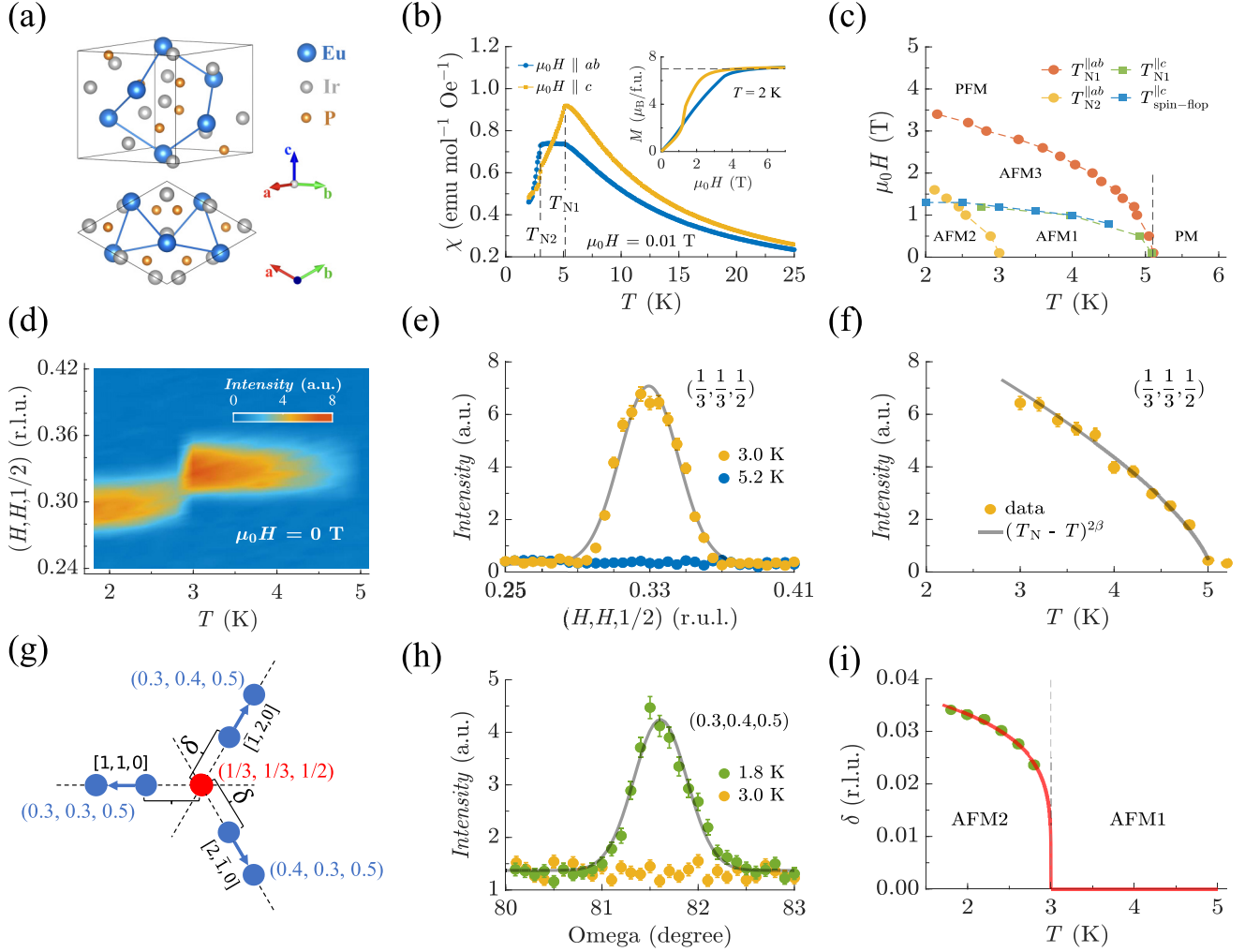


FIG. 1. Crystalline structure and magnetic properties of EuIr_2P_2 . (a) Schematic of crystalline structure of EuIr_2P_2 . (b) Magnetic susceptibilities for $\mu_0 H \parallel ab$ and $\mu_0 H \parallel c$ with a magnetic field $\mu H = 0.01$ T, respectively. The inset is the magnetization as a function of magnetic field at $T = 2$ K. (c) Extracted $\mu_0 H - T$ phase diagram. (d) Single-crystal neutron-diffraction measurement of the temperature dependence of the magnetic satellite peak, shown as a contour map of $(H, H, 1/2)$ reciprocal-space scans vs temperature. (e) $(H, H, 1/2)$ scans around $(1/3, 1/3, 1/2)$ at $T = 3$ and 5.2 K, respectively. The gray line is the Gaussian fitting. (f) Temperature dependence of the integrated intensity of $(1/3, 1/3, 1/2)$ and the gray line is the fitting curve by using exponential decay function. (g) A schematic representation in reciprocal space that illustrates the transformation of the magnetic satellite peak $(1/3, 1/3, 1/2)$ in AFM1 to the incommensurate magnetic peaks that emerged in AFM2. (h) Sample-rotation scans (i.e., Omega scans) at $(0.3, 0.4, 0.5)$ at $T = 1.8$ and 3.0 K, respectively. The gray line is the Gaussian fitting. (i) Temperature dependence of the incommensurability δ , determined from the magnetic satellite peak $(1/3 - \delta, 1/3 - \delta, 1/2)$. The solid red line is a guide to the eye.

which were consistent with theoretical value of Eu^{2+} ($\mu_{\text{sat}} = gJ\mu_B$, where $g = 2$ is the Landé factor for Eu^{2+}) wherein the orbital magnetic moment is completely quenched. The magnetic field-temperature ($\mu_0 H - T$) phase diagram summarized in Fig. 1(c) is based on the data extracted from magnetic properties (see Fig. S3 in SM) [24].

To gain insight into the magnetic structures within the various ordered phases, we carried out a detailed single-crystal neutron diffraction (SCND) experiment on one large piece of EuIr_2P_2 single crystal (see Fig. S4a in SM). Figure 1(d) shows the temperature dependence of the $(H, H, 1/2)$ scans in reciprocal space under zero magnetic field, revealing two transitions at around $T_{N1} = 5.0$ K and $T_{N2} = 2.9$ K,

respectively, in consistence with the magnetic susceptibility measurements. The first transition at T_{N1} is from the paramagnetic (PM) phase to a commensurate $\sqrt{3} \times \sqrt{3} \times 2$ AFM phase (also referred to 120° AFM phase), namely AFM1, with the magnetic propagation vector of $k_{\text{AFM1}} = (1/3, 1/3, 1/2)$. This transition is clearly illustrated in Fig. 1(e), where the magnetic satellite peak $(1/3, 1/3, 1/2)$ observed at 3.0 K completely disappears above $T_{N1} \sim 5.0$ K. Additionally, the temperature dependence of the integrated intensity of $(1/3, 1/3, 1/2)$ agrees well with the fitting from the exponential decay function as shown in Fig. 1(f). The estimated exponent $\beta = 0.34$ implies a three-dimensional Heisenberg character for the magnetism in this compound [26]. With further

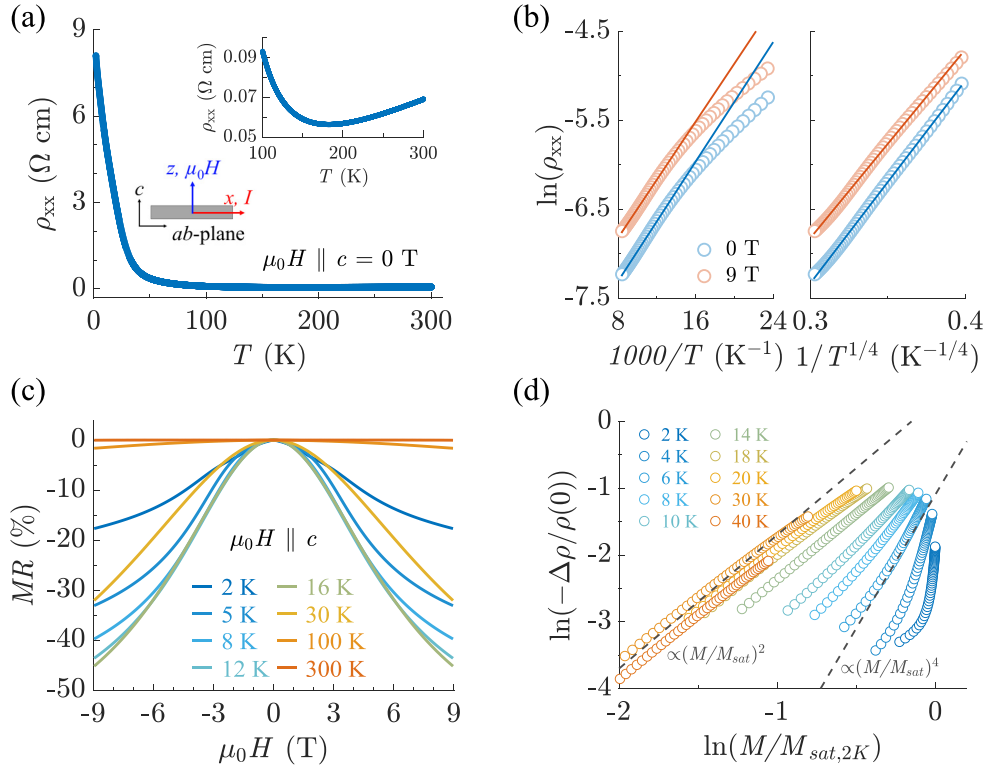


FIG. 2. Longitudinal resistivity and magnetoresistance of EuIr_2P_2 . (a) Longitudinal resistivity as a function of temperature when current is injected in the ab plane and magnetic field is applied along c axis. The lower inset is a schematic of geometry for the transport measurement and the upper inset is enlarged view in temperature range from 100 to 300 K. (b) Arrhenius- and Mott-law fitting of the temperature-dependent longitudinal resistivity in temperature range from 40 to 120 K, respectively. (c) Magnetoresistance as functions of magnetic field at various temperatures. (d) Logarithmic plot of the normalized magnetization-dependent magnetoresistance at various temperatures.

decreasing temperature to T_{N2} , an abrupt phase transition from the commensurate $\sqrt{3} \times \sqrt{3} \times 2$ AFM phase to an incommensurate AFM phase, namely AFM2, takes place at T_{N2} . This transition is signalized by the sudden disappearance of $(1/3, 1/3, 1/2)$ and the concomitant emergence of more incommensurate magnetic satellite peaks, for instance, $(0.3, 0.3, 0.5)$, and $(0.3, 0.4, 0.5)$ and $(0.4, 0.3, 0.5)$, etc., seen at 1.8 K [see Fig. 1(g) and Fig. 1(h)]. These incommensurate magnetic peaks are also observed around other symmetry-equivalent Q positions of $(1/3, 1/3, 1/2)$, such as around $(-1/3, 2/3, 1/2)$ and $(2/3, 2/3, 1/2)$, etc. The magnetic propagation vector of the AFM2 phase can thus be identified as $k_{\text{AFM2}} = (1/3 \pm \delta, 1/3 \pm \delta, 1/2)$ (or the symmetry-equivalent $[1/3 \mp \delta, 1/3 \pm \delta, 1/2]$ and $(1/3 \pm \delta, 1/3 \mp \delta, 1/2)$), where the incommensurability $\delta \approx 0.0341$ r.l.u. at 1.8 K. Below T_{N2} , the incommensurability δ gradually increases during cooling, as shown in Fig. 1(i). The first-order-like magnetic phase transition from AFM1 to AFM2 can also be clearly seen from the abrupt drop in magnetic susceptibilities for both $\mu_0 H \parallel ab$ and $\mu_0 H \parallel c$ at T_{N2} [see Fig. 1(b)]. The determination of the magnetic structures in both AFM1 and AFM2 phases via neutron diffraction is particularly challenging because of the strong neutron absorption by the Eu and Ir elements in this material and the high complexity of the magnetic structure itself. Nevertheless, based on the various simulated magnetic structures of AFM1 phase [see Figs. S5(a)–S5(c) in SM], it turns out that the Eu magnetic moments, or the projections of the Eu magnetic moments within the ab plane, likely form a

commensurate 120° clamped AFM structure. This is in contrast to the previous theoretical models that suggest a collinear AFM structure in this temperature range [23]. With a finite magnetic field applied along the c axis, a noncoplanar spin texture could nevertheless be realized even in the AFM1 phase [see Fig. S5(d) in SM]. Furthermore, based on the determined magnetic propagation vector k_{AFM2} and magnetic susceptibilities, a long-periodicity incommensurate spiral magnetic order, with a spiral pitch length at ~ 98 Å at 1.8 K, appears a plausible scenario for AFM2. Such type of helical magnetic order can be an ideal fertile ground for the emergence of topologically nontrivial noncoplanar spin textures under applied magnetic fields, like magnetic skyrmions in MnSi [26]. While our neutron-diffraction results and analysis strongly hint that EuIr_2P_2 could be a candidate of AFM skyrmion compounds, more direct evidence from further investigations should be addressed in the future. Having addressed the magnetic structures in EuIr_2P_2 , we will further demonstrate the large THE in this compound.

We then show the magnetoelectric transport properties of this compound. To perform the measurements, the crystal was cut into a thin plate with the ab plane exposed, allowing the magnetic field to be applied along the c axis as shown in the inset of Fig. 2(a). To avoid the possible current-activated carrier, a small electric current with low frequency according to the current (I)-dependent voltage (V) measurement is adopted as shown in Fig. S6(b) in SM. The temperature-dependent longitudinal resistivity ρ_{xx} , as plotted in Fig. 2(a), exhibits

semiconducting behavior with decreasing temperature. In the enlarged temperature range between 100 and 300 K, as shown in the inset of Fig. 2(a), a metallic behavior is observed when the temperature is reduced to about 180 K, and then the semiconducting behavior is observed on further cooling. This suggests the impurity- or vacancy-induced energy levels in the band gap. To understand the semiconducting behavior below 180 K, an Arrhenius model ($\rho_{xx} \sim e^{E_g/k_B T}$, E_g is energy gap and k_B is Boltzmann constant) was first utilized to fit the observed data in temperature range from 40 to 120 K, as shown in the left panel of Fig. 2(b). The estimated energy gap is around 14 meV, which is largely reduced in contrast to the theoretical value of 400 meV [22]. The estimated energy gaps at different magnetic fields are also presented in Figs. S6(c) and S6(d) in SM, showing a weak magnetic field dependence. The divergence between observed data and fitted line suggests that the thermal activation cannot completely describe the longitudinal resistivity in this compound. This behavior recalls that in extrinsic semiconductors, because of the typical resistances between neighboring impurities, become larger than those connecting some distant impurities whose energy levels happen to be very close to the Fermi level [27]. A Mott variable range-hopping (VRH) model ($\rho_{xx} \propto e^{(T_M/T)^{1/4}}$; T_M is Mott temperature) is therefore required to describe the temperature-dependent resistivity [28,29]. As a result, the resistivity data in temperature range from 80 to 120 K can be well reproduced by using Mott's VRH law as shown in the right panel of Fig. 2(b), yielding $T_M \sim 27000$ and $20\,000$ K for $\mu_0 H = 0$ and 9 T, respectively. These values are close to the VRH oxide insulator [30]. The hopping carrier not only dominates the longitudinal resistivity here, but also significantly dominates the observed THE in this compound, which we will discuss later. At lower temperature close to the magnetic phase transition, as plotted in Fig. S6(b) in SM, the resistivity decreases with increasing magnetic field, performing a negative magnetoresistance ($MR = \frac{\rho_H - \rho_{H=0}}{\rho_{H=0}} \times 100\%$). The MRs as a function of magnetic field for various temperatures are plotted in Fig. 2(c). Interestingly, the largest negative MRs occur around 16 K, which is larger than three times the T_{N1} , suggesting a spin-fluctuation state in the paramagnetic (PM) phase. To gain insight into the relationship between MR and magnetic properties in EuIr_2P_2 , we plot the MR as a function of the normalized magnetization ($M/M_{\text{sat}, 2\text{ K}}$) in logarithmic scale as shown in Fig. 2(d). At temperature above 16 K, all curves follow a linear relationship with magnetization formulated by $MR \propto (M/M_{\text{sat}, 2\text{ K}})^2$. This behavior is typically observed in colossal MR (CMR) systems [31,32]. As the temperature is further reduced, the observed MR follows $MR \propto (M/M_{\text{sat}, 2\text{ K}})^4$, which is similar to the result of spin fluctuation as in weakly ferromagnetic metals [33]. When the temperature is in the AFM phase, the exponent becomes even larger. According to the band calculation (see Fig. S9 in SM), the Eu $4f$ bands are located approximately 1 eV below the Fermi level, but the local moments still hybridize with the hopping carriers. The negative MR above T_{N1} can be attributed to the possible spin-fluctuation mechanism as previously reported in CMR systems. Below T_{N1} , the scattering from local moments will be suppressed when the moments are forcibly reoriented along the direction of magnetic field, resulting in a negative MR. However, the high exponents extracted from

the scaling relation between MR and magnetization in the AFM phase require further investigation.

The analysis above shows that the hopping carrier is crucial and strongly linked to the magnetic structure in EuIr_2P_2 . Then, we focused on the observation of the anomalous Hall effect (AHE) and THE in this compound. In time-reversal symmetry (TRS) broken systems, such as FM compounds, the spontaneous magnetization induces a nonzero Hall resistivity at zero magnetic field, which is known as AHE. Meanwhile, the zero-field AHE occurs also in AFM materials when the magnetic structure becomes noncollinear, such as Mn_3Sn and Mn_3Ge [34,35]. The AHE is also observed in soft ferromagnets and collinear AFM materials, and is typically proportional to the magnetization. In this case, there is no zero-field AHE but a magnetization-dependent Hall resistivity that will saturate when the moments are fully polarized by external magnetic field. In this paper we show that in EuIr_2P_2 , the AHE is proportional to the high-field magnetization and nonproportional to the low-field magnetization. The isothermal magnetization of EuIr_2P_2 for various temperatures is shown in Fig. 3(a), where a spin-flop transition occurred at approximately $\mu_0 H \sim 1.3$ T below 4 K. As shown in Fig. 3(b), the total Hall resistivity as a function of magnetic field exhibited nonlinear behavior below 20 K. Besides, an obvious saturation of Hall resistivity occurs at $\mu_0 H \sim 3$ T that is consistent with the saturated magnetic field of magnetization, implying the AHE in EuIr_2P_2 . The observed Hall resistivity of magnetic compound is usually described as [18,36]

$$\rho_{yx} = \rho_{yx}^N + \rho_{yx}^A + \rho_{yx}^T = R_0 \mu_0 H + S_H \rho_{xx}^n M + \rho_{yx}^T, \quad (1)$$

where $\rho_{yx}^N = R_0 \mu_0 H$ is ordinary Hall effect (OHE), which is perpendicular to the allied magnetic field, $\rho_{yx}^A = S_H \rho_{xx}^n M$ is conventional AHE, which is related to magnetization and n is 1 or 2 depending on scattering mechanism by spontaneous magnetization [9], and ρ_{yx}^T is the THE, which is considered to be a characteristic of topological properties arising from electronic bands or other magnetic structure. According to this process, the holelike carrier with a concentration of $\sim 10^{19} \text{ cm}^{-3}$ and mobility of $\sim 1 \text{ cm}^2 \text{ V}^{-1} \text{ s}^{-1}$ can be estimated [see Fig. S8(a) in SM].

Using Eq. (1), the OHE, AHE, and THE can be extracted from the observed Hall data with the assistance of magnetization. The focus is on the AHE behavior at high magnetic field range. By fitting the Hall resistivity in the high-field range via a linear function, the intercept value is the AHE resistivity ρ_{yx}^A [also see Fig. S7(b) in SM]. To understand the mechanism of AHE observed in EuIr_2P_2 , ρ_{yx}^A is plotted as a function of zero-field ρ_{xx} in Fig. 3(c). A linear scaling relation has been identified. In good conductors ($\sigma_{xx} > 10^6 \Omega^{-1} \text{ cm}^{-1}$), the linear relation between ρ_{yx}^A and ρ_{xx} indicates the skew-scattering mechanism [9]. However, in EuIr_2P_2 , the conductivity σ_{xx} is several orders of magnitude smaller ($\sigma_{xx} < 1 \Omega^{-1} \text{ cm}^{-1}$), which cannot be well understood with the same mechanism. The low conductivity also places EuIr_2P_2 into the hopping carrier regime, as shown in Fig. 3(d). The extracted anomalous Hall conductivity (AHC, $\sigma_{xy}^A = \frac{\rho_{yx}^A}{\rho_{xx}^2 + \rho_{yx}^2}$) follows the empirical power-law relationship with the longitudinal conductivity ($\sigma_{xy}^A \propto \sigma_{xx}^{1.6}$), indicating

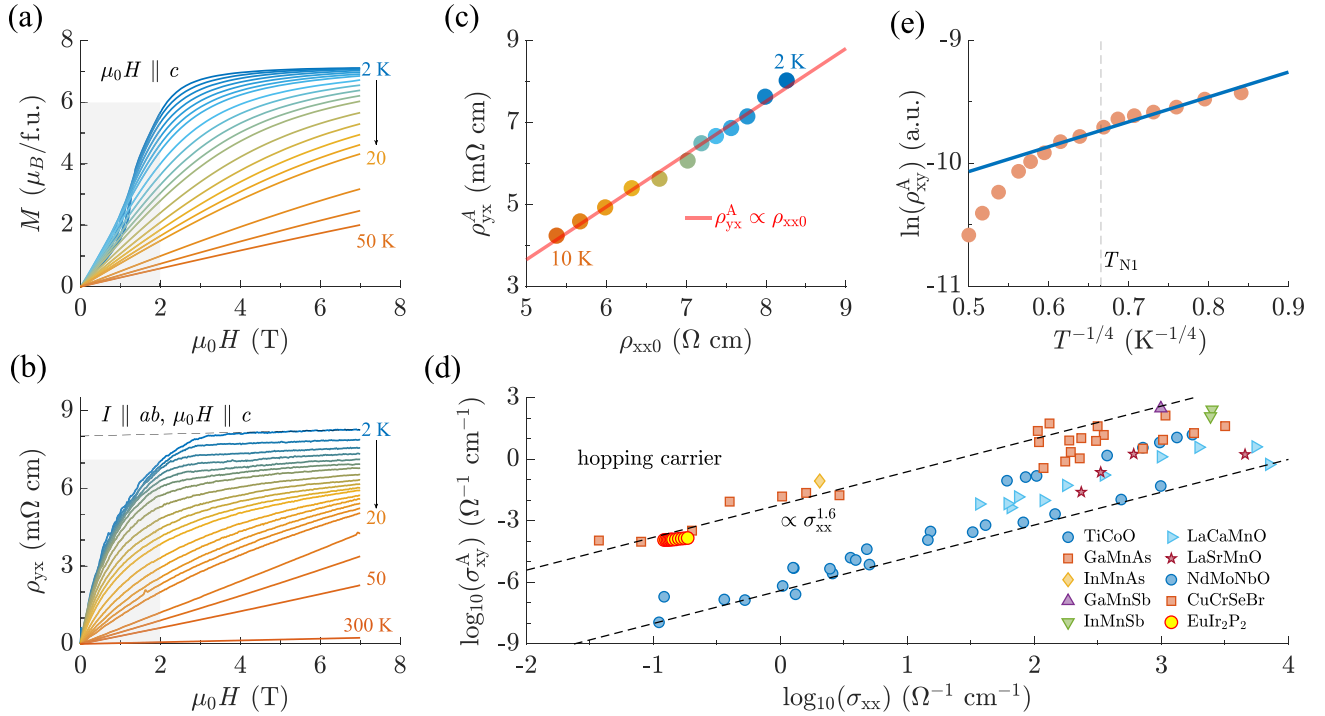


FIG. 3. Anomalous Hall effect (AHE) below and above T_{N1} . (a) The isothermal magnetization and (b) total Hall resistivity as functions of magnetic field, respectively. Temperature gradient: every 0.5 K one step for 2 – 5 K, every 1 K one step for 5 – 10 K, every 2 K one step for 10 – 20 K, and every 10 K one step for 20 – 50 K. Gray shadows in (a) and (b) highlight the difference between magnetization and Hall resistivity. (c) Anomalous Hall resistivity as a function of zero-field longitudinal resistivity at various temperatures from 2 to 10 K. Red solid line is the linear fitting of $\rho_{yx}^A \propto \rho_{xx}$. (d) σ_{xy}^A as a function of zero-field σ_{xx} for EuIr_2P_2 and other selected compounds in hopping regime. The data are extracted from previous reports [4,51–60]. Dashed gray lines are guides for the empirical relationship of $\sigma_{xy}^A \propto \sigma_{xx}^{1.6}$. (e) Temperature-dependent anomalous Hall resistivity fitted by the model $\rho_{yx}^A \propto e^{(T_M/T)^{1/4}}$ (blue line).

the extrinsic contribution of anomalous Hall effect (AHE) in the hopping regime of this compound. Meanwhile, given the semiconducting nature of EuIr_2P_2 , it is necessary to exclude any subjective evaluations of the intrinsic contribution of the AHE that may arise from the BC of the electronic bands. The AHE observed here is then compared with that of some representative magnetic semiconductors. It is worth noting that in some Eu-based compounds, such as EuAs thin film [37] and EuTe [38], a spin cluster skew-scattering mechanism is adopted to explain the observed AHE. In this scenario, the AHE is nonproportional to the magnetization and originates from the thermally fluctuating spins that form the local noncoplanar spin structures, even far above the Néel temperature. However, the AHE observed in EuIr_2P_2 is different. Upon comparing Figs. 3(a) and 3(b) [also see Figs. S7(a) and S7(b) in SM], it is evident that the magnetization and Hall resistivity are saturated at high magnetic fields and below T_{N1} , where the magnetic moments are fully polarized and form a net magnetization. The divergence between magnetization and Hall resistivity [see Fig. S7(a) in SM] at low magnetic field is the contribution of THE that will be discussed later. According to a theory proposed by Burkov and Balents [39], the microscopic origin of the AHE in diluted ferromagnetic semiconducting systems, such as $\text{Ga}_{1-x}\text{Mn}_x\text{As}$ [40], can be analogized to the geometrical Berry phase mechanism in metallic ferromagnets. It suggested that the carrier gains a phase when hopping around closed-loop paths in the presence

of spin-orbit interactions (SOIs) and the background magnetization of the localized magnetic moments. In EuIr_2P_2 , the localized moments of Eu^{2+} are largely correlated with the hopping conducting carrier and also provide a large SOI. This implies a similar mechanism to this theoretical proposal. According to this theory, a quantitative relation, formulated by $\rho_{yx}^A \propto e^{(T_M/T)^{1/4}}$, is adopted to analyze the observed data and gives $T_M \sim 25$ K, as shown in Fig. 3(e). The observed anomalous Hall resistivity fits well to Mott's law below 10 K but diverges at high temperature. Due to the limitations of the magnetic field during measurement, the extracted value of ρ_{yx}^A may be relatively underestimated as a result of non-saturated magnetization at higher temperature above T_{N1} . Meanwhile, the estimated T_M appears to be relatively small. This can be attributed to the reduced effectiveness of Mott's VRH mechanism under high magnetic fields [41]. Therefore, the T_M estimated from the temperature-dependent anomalous Hall resistivity, which is deduced from high-field data, is underestimated.

We now turn our attention to the THE observed in the low magnetic field range in EuIr_2P_2 . Our investigation, based on neutron-scattering data, reveals that EuIr_2P_2 undergoes a commensurate noncollinear or noncoplanar magnetic structure in the AFM1 phase ($T_{N2} < T < T_{N1}$), and a more intricate incommensurate magnetic structure in the AFM2 phase ($T < T_{N2}$). In the presence of such a magnetic structure, THE is usually induced by the real-space BC when the three nearest

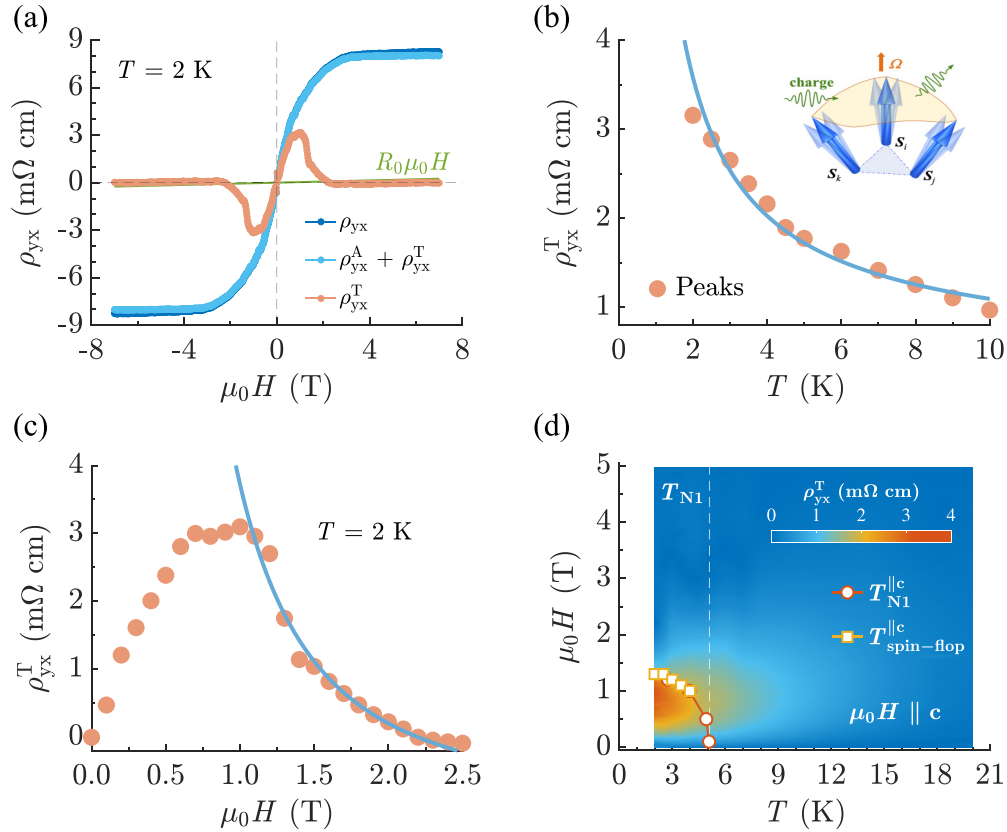


FIG. 4. Mechanism of topological Hall effect (THE). (a) Magnetic field-dependent total Hall resistivity, extracted ordinary Hall resistivity, anomalous Hall resistivity, and topological resistivity at $T = 2$ K, respectively. (b) Fitting of the topological resistivity by Mott's variable range-hopping law (the solid blue line). The inset is a schematic to illustrate the formation of spin chirality and the generated Berry curvature. (c) Magnetic field-dependent topological Hall resistivity. The solid line is the fitting curve by $\rho_{yx}^T \propto 1/(\mu_0H)^2$. (d) Temperature-magnetic field phase diagram of topological Hall resistivity and magnetic phase transitions.

spins form a spin chirality structure, as documented in previous studies [1,2,4,5]. It is noteworthy that THE is expected to diminish at high magnetic fields due to the coalignment of magnetic moments [36]. Consequently, Eq. (1) can be reformulated as follows:

$$\rho_{yx}/\mu_0H = R_0 + S_H \rho_{xx}^n M/\mu_0H. \quad (3)$$

Following this process, the THE can be extracted by subtracting the OHE and magnetization-proportional AHE [see Figs. S7(c) and S7(d) in SM]. Figure 4(a) shows the separating results, which illustrate each contribution to the total Hall resistivity. Normally, the MR is ideally sufficiently low for convincingly considering the value of $n = 1$ or 2 according to different scattering mechanism, where $n = 1$ implies the occurrence of skew-scattering mechanism and $n = 2$ indicates that the side jump or intrinsic scattering mechanism is dominating [9,42]. Meanwhile, $n = 0$ implies scaling of the AHE by magnetization via the simple relationship of $\rho_{yx}^A = R_S 4\pi M$. In contrast, the separation processes using parameter $n = 0, 1, 2$ exhibited negligible difference in the hump feature at $\mu_0H \sim 1$ T no more than a declination at $\mu_0H \sim 2.5$ T [see Fig S7(d) in SM]. Considering the relationship of $\rho_{yx}^A \propto \rho_{xx}$ as shown in Fig. 3(c), $n = 1$ was adopted during the procedure of separation for Hall resistivity. As a result, the yielded value of ρ_{yx}^T reaches around 3 m Ω cm at 2 K and 1 m Ω cm at

10 K. This value exceeds several orders of magnitude beyond chiral (achiral) crystalline materials that exhibit THE. The observed topological Hall resistivity is compared with other representative materials as listed in Table I. EuIr₂P₂ displays the highest value, surpassing all reported crystalline materials. The large ρ_{yx}^T of EuIr₂P₂ profits from its spin chirality induced by field as well as its semiconducting nature. Generally, the reported THE can be separated into two regimentations: (i) Spin chirality structure-dominated systems such as MnSi [1], FeGe [5], Gd₂PdSi₃ [2], etc., which refers to the formation of skyrmions. (ii) Weyl point-dominated systems such as EuCd₂As₂ [18], GdPtBi [14,15], and EuB₆ [19], in which the THE is presented and expounded via topological band structures with nonzero BC in momentum space. As EuIr₂P₂ is a semiconductor with a relative energy gap, as well as the observed noncollinear magnetic structure in EuIr₂P₂ (see Fig. S5 in SM), the spin chirality plays a decisive role for the experimental discovery of THE. Below T_{N1} , when applying the magnetic field along the c axis, the Eu²⁺ moments will be forcedly reorientated along the c axis. The three nearest Eu²⁺ moments will form a noncoplanar spin chirality structure, thus giving a nonzero real-space BC, as schematized in the inset of Fig. 4(b). Above T_{N1} , the Eu²⁺ moments on triangular lattice are also expected to form the noncoplanar fluctuation with finite spin chirality [37], contributing to the THE.

TABLE 1. Summary of the topological Hall resistivity of current studied crystalline materials.

Materials	$\rho_{xx}(\mu\Omega \text{ cm})$	$\rho_{yx}^A(\mu\Omega \text{ cm})$	$\rho_{yx}^T(\mu\Omega \text{ cm})$	Reference
MnSi ^a	0.4–5	0.1	0.0045	[1]
EuPtSi ^a	0.025	0.5	0.12	[3]
Gd ₂ PdSi ₃ ^b	70		2.5	[2]
EuIn ₂ As ₂ ^b	120	2	1	[12]
EuCuAs ^b	150	2.4	7.4	[13]
EuCd ₂ As ₂ ^b	4000	10	13	[18]
MnBi ₄ Te ₇ ^b	1000	12	7	[20]
EuIr ₂ P ₂ ^a (@ 2 K)	8.1×10^6	8028	3156	This work
EuIr ₂ P ₂ ^a (@ 10 K)	5.2×10^6	4245	970	This work

^aChiral crystalline structure.^bAchiral crystalline structure.

Owing to Mott's VRH carrier in this compound, we proposed a qualitative theory based on Mott's VRH law to understand the temperature- and magnetic field dependences of THE (see Appendix A for details). A possible mechanism could involve temperature-dependent phonon-mediated carrier hopping between different sites which were claimed to be the case for a similar observation performed for EuAs [37]. However, this consideration was given to large temperatures $T \sim 200$ K and the leading temperature dependence was thus provided by a temperature-dependent phonon density. Since the $\rho_{yx}(T, B)$ measurements are conducted well below Debye temperature $T_D \sim 280$ K [23], we suggest that the observed THE has a pure disorder origin, and the temperature dependence itself is given by Mott's exponent. As shown in Fig. 4(b), the peak values ρ_{yx}^T as a function of temperature can be well reproduced by Mott's VRH law, expressed by $\rho_{yx}^T \propto e^{(T_M/T)^{1/4}}$, giving the Mott temperature $T_M \sim 327$ K. According to our theoretical modeling, the magnetic field-dependent ρ_{yx}^T follows a $1/(\mu_0 H)^2$ decaying due to the existence of spin chirality as shown in Fig. 4(c). The $\mu_0 H - T$ phase diagram is then plotted in Fig. 4(d), illustrating an area of the observed ρ_{yx}^T . Similar phase diagrams have also been observed in many skyrmion materials, including Gd₂PdSi₃ [2], EuPtSi [3], and FeGe [5]. This suggests that EuIr₂P₂ is an intriguing candidate for an AFM skyrmion material.

IV. CONCLUSION

In summary, we report the observation of a large THE in the AFM semiconductor EuIr₂P₂ with a chiral crystal structure. The experimentally observed spiral magnetic ground state is attributed to the unique chiral symmetry of this material. The VRH carriers interact significantly with the spin chirality induced by an external magnetic field, resulting in the emergence of THE. Due to the ultralow conductivity of EuIr₂P₂, the THE is predominantly governed by spin chirality-induced Berry curvature. This characteristic renders EuIr₂P₂ a rare and exceptional case that fills the gap in THE research within the ultralow-conductivity hopping regime. Our findings expand the scope of research in this transport regime and present opportunities for future studies. However, further experimental investigation is required to provide direct evidence of skyrmions. Additionally, the semiconducting

nature of EuIr₂P₂ offers a promising platform for exploring and applying field-tunable devices and semiconducting AFM spintronics.

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APPENDIX A: THEORETICAL MODEL OF TOPOLOGICAL HALL EFFECT

The problem of anomalous Hall conductivity (AHC) in hopping regime was previously theoretically studied with the help of the following Hamiltonian [39]:

$$\begin{aligned}
 H &= H_e + H_{ph} + H_{e-ph}, \quad H_e \\
 &= \sum_{i\alpha} \varepsilon_i c_{i\alpha}^\dagger c_{i\beta} - \sum_{ij\alpha} t_{ij\alpha} c_{i\alpha}^\dagger c_{j\alpha} - \sum_{i\alpha\beta} h S_{\alpha\beta} c_{i\alpha}^\dagger c_{i\beta}, \quad H_{ph} \\
 &= \sum_k \omega_k b_k^\dagger b_k, \quad H_{e-ph} = g \sum_{k,p,\alpha} c_{p+k,\alpha}^\dagger c_{p,\alpha} (b_k + b_{-k}^\dagger),
 \end{aligned} \tag{A1}$$

where i stands for the ion number, α describes spin degrees of freedom, and h is a TRS breaking factor (for instance, it could describe a nonvanishing net magnetic moment proportional to h). For Hamiltonian (4) it was proven that AHC demonstrates Mott's temperature dependence. However, consideration was given only to the system with explicitly broken TRS [39]. For noncoplanar spin configurations it was proven that hopping amplitude has a nontrivial spin chirality-related phase factor: $t_{i\alpha j\alpha} = t_{\alpha} t_{ij} \propto t_{\alpha} e^{i\phi_{ij} - r_{ij}/\xi}$ [43]. Here, r_{ij} and ξ are interatomic distance and localization length, correspondingly, and ϕ_{ij} is a

direct analog of Berry phase $\phi_{ij} = \int_{r_i}^{r_j} A dr$, which plays the role of TRS breaking factor instead of h . Thus, we arrive at the following result for anomalous Hall response j_{AH} that in insulating regime the Hall response is cubic in t [39,44]:

$$j_{\text{AH}, i \rightarrow j} \propto E \sum_m \text{Im}\{t_{ij} t_{jm} t_{mi}\} \propto E \langle \phi_{ijk} \rangle e^{-(T_M/T)^{1/4}}, \quad (\text{A2})$$

where $\langle \phi_{ijk} \rangle \equiv \langle S_i \cdot \langle S_j \times S_k \rangle \rangle$ stands for averaged spin chirality [9] and Mott exponent $e^{-(T_M/T)^{1/4}}$ describes the scattering rate T dependence as shown in Fig. 4(b). However, one should be aware that the result of Eq. (A2) is applicable only to the case of small net momentum, i.e., for $T < T_N$. Above the T_N anomalous Hall current would be given not only by $\langle \phi_{ijk} \rangle$ but rather by some mix of finite magnetization and average spin chirality contributions. Temperature dependence of spin

chirality factor could be neglected since exchange interaction was previously estimated to be small (~ 0.5 K) [22].

As shown in the main text, the most striking difference between ρ_{yx}^T and ρ_{yx}^A lies in different magnetic field dependence (nonmonotonic for ρ_{yx}^T vs saturating for ρ_{yx}^A). This observation implies that $\rho_{yx}^T(B_{\text{ext}})$ behavior should have its origin in magnetic and not disorder properties (which is basically the same for both contributions). Topological Hall response behavior with B_{ext} is, generally speaking, heavily dependent on spin interaction details in EuIr_2P_2 which so far remain unknown. Nonetheless, we provide a simple theoretical model that allows for decaying behavior of $\rho_{yx}^T(B_{\text{ext}})$. Given that the spin structure is noncoplanar, minimal ingredients for the model are Zeeman term and spin interaction term that favors noncoplanar structure (e.g., Dzyaloshinskii-Moriya term is known to stabilize out-of-plane spin structures). Using the results obtained in the previous section, for $j_{\text{AH}}(B_{\text{ext}}) \sim \langle \phi_{ijk} \rangle(B_{\text{ext}})$ we have

$$\langle \phi_{ijk} \rangle = \langle S_i \cdot (S_j \times S_k) \rangle = \frac{\int DSS_i \cdot (S_j \times S_k) e^{-\beta \mu_B \sum_p B_{\text{ext}} \cdot S_p - \beta \sum_{p,q} D(S_p \times S_q)}}{\int DSe^{-\beta \sum_p B_{\text{ext}} \cdot S_p - \beta \sum_{p,q} D_{p,q} \cdot (S_p \times S_q)}}, \quad (\text{A3})$$

where $\beta = 1/(k_B T)$ and the integral term describes angle average over classical spin direction. In the limit of $\mu_B B_{\text{ext}} \gg D$ we arrive at

$$\langle \phi_{ijk} \rangle \approx \frac{\int DSS_i \cdot (S_j \times S_k) (-\beta \sum_{p,q} D_{p,q} \cdot (S_p \times S_q)) e^{-\beta \mu_B \sum_p B_{\text{ext}} \cdot S_p}}{\int DSe^{-\beta \mu_B \sum_p B_{\text{ext}} \cdot S_p}}. \quad (\text{A4})$$

Switching to dimensionless units $\beta \mu_B B_{\text{ext}} \rightarrow B$, $\beta D \rightarrow D$, and evaluating the integral, we find

$$\langle \phi_{ijk} \rangle \approx \frac{D \cdot \hat{B}}{B^2}, \quad \hat{B} = \frac{B}{|B|}. \quad (\text{A5})$$

The fitting of ρ_{yx}^T as a function of B_{ext} with Eq. (A5) is presented at Fig. 4(c). Note that the B_{ext} is $\mu_0 H$ in the main text.

APPENDIX B: MAGNETIC STRUCTURE SIMULATION AND BAND-STRUCTURE CALCULATION

The simulated magnetic structures (see Fig. S5 in SM) in the AFM1 phase are performed using the MAXMAGEN [45]. The electronic structure calculations were performed based

on the density-functional theory (DFT) using the plane-wave projected augmented-wave method as implemented in Vienna *Ab initio* Simulation Package (VASP) [46–48]. A plane-wave cutoff of 500 eV and a k -point mesh of $9 \times 9 \times 7$ were used for the self-consistent calculations. We used the Perdew-Burke-Ernzerhof exchange-correlation functional within the generalized gradient approximation [49]. The strongly localized Eu 4*f* electrons were treated using the DFT + U method [50], where the onsite Coulomb interaction U is considered to be 5 eV.

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- [1] A. Neubauer, C. Pfleiderer, B. Binz, A. Rosch, R. Ritz, P. G. Niklowitz, and P. Böni, Topological Hall effect in the A phase of MnSi, *Phys. Rev. Lett.* **102**, 186602 (2009).
 - [2] T. Kurumaji, T. Nakajima, M. Hirschberger, A. Kikkawa, Y. Yamasaki, H. Sagayama, H. Nakao, Y. Taguchi, T. Arima, and Y. Tokura, Skyrmion lattice with a giant topological Hall effect in a frustrated triangular-lattice magnet, *Science* **365**, 914 (2019).
 - [3] M. Kakihana *et al.*, Giant Hall resistivity and magnetoresistance in cubic chiral antiferromagnet EuPtSi, *J. Phys. Soc. Jpn.* **87**, 023701 (2018).
 - [4] Y. Taguchi, Y. Oohara, H. Yoshizawa, N. Nagaosa, and Y. Tokura, Spin chirality, Berry phase, and anomalous Hall effect in a frustrated ferromagnet, *Science* **291**, 2573 (2001).
 - [5] E. Turgut, M. J. Stolt, S. Jin, and G. D. Fuchs, Topological spin dynamics in cubic FeGe near room temperature, *J. Appl. Phys.* **122**, 183902 (2017).
 - [6] Y. Fujishiro *et al.*, Topological transitions among skyrmion- and hedgehog-lattice states in cubic chiral magnets, *Nat. Commun.* **10**, 1059 (2019).
 - [7] N. Kanazawa, Y. Onose, T. Arima, D. Okuyama, K. Ohoyama, S. Wakimoto, K. Kakurai, S. Ishiwata, and

- Y. Tokura, Large topological Hall effect in a short-period helimagnet MnGe, *Phys. Rev. Lett.* **106**, 156603 (2011).
- [8] H. Wang, Y. Dai, G.-M. Chow, and J. Chen, Topological Hall transport: Materials, mechanisms and potential applications, *Prog. Mater. Sci.* **130**, 100971 (2022).
- [9] N. Nagaosa, J. Sinova, S. Onoda, A. H. MacDonald, and N. P. Ong, Anomalous Hall effect, *Rev. Mod. Phys.* **82**, 1539 (2010).
- [10] H. Takagi *et al.*, Spontaneous topological Hall effect induced by non-coplanar antiferromagnetic order in intercalated van der Waals materials, *Nat. Phys.* **19**, 961 (2023).
- [11] C. Sürgers, G. Fischer, P. Winkel, and H. V. Löhneysen, Large topological Hall effect in the non-collinear phase of an antiferromagnet, *Nat. Commun.* **5**, 3400 (2014).
- [12] J. Yan, Z. Z. Jiang, R. C. Xiao, W. J. Lu, W. H. Song, X. B. Zhu, X. Luo, Y. P. Sun, and M. Yamashita, Field-induced topological Hall effect in antiferromagnetic axion insulator candidate EuIn₂As₂, *Phys. Rev. Res.* **4**, 013163 (2022).
- [13] S. Roychowdhury *et al.*, Interplay between magnetism and Topology: Large topological Hall effect in an antiferromagnetic topological insulator, EuCuAs, *J. Am. Chem. Soc.* **145**, 12920 (2023).
- [14] T. Suzuki, R. Chisnell, A. Devarakonda, Y.-T. Liu, W. Feng, D. Xiao, J. W. Lynn, and J. G. Checkelsky, Large anomalous Hall effect in a half-Heusler antiferromagnet, *Nat. Phys.* **12**, 1119 (2016).
- [15] C. Shekhar *et al.*, Anomalous Hall effect in Weyl semimetal half-Heusler compounds RPtBi (R = Gd and Nd), *Proc. Natl. Acad. Sci. USA* **115**, 9140 (2018).
- [16] T. Shang, Y. Xu, D. J. Gawryluk, J. Z. Ma, T. Shiroka, M. Shi, and E. Pomjakushina, Anomalous Hall resistivity and possible topological Hall effect in the EuAl₄ antiferromagnet, *Phys. Rev. B* **103**, L020405 (2021).
- [17] M. M. Piva, J. C. Souza, G. A. Lombardi, K. R. Pakuszewski, C. Adriano, P. G. Pagliuso, and M. Nicklas, Topological Hall effect in CeAlGe, *Phys. Rev. Mater.* **7**, 074204 (2023).
- [18] Y. Xu *et al.*, Unconventional transverse transport above and below the magnetic transition temperature in Weyl semimetal EuCd₂As₂, *Phys. Rev. Lett.* **126**, 076602 (2021).
- [19] J. Shen *et al.*, Magnetic-field modulation of topological electronic state and emergent magneto-transport in a magnetic Weyl semimetal, *The Innovation* **4**, 100399 (2023).
- [20] S. Roychowdhury, S. Singh, S. N. Guin, N. Kumar, T. Chakraborty, W. Schnelle, H. Borrmann, C. Shekhar, and C. Felser, Giant topological Hall effect in the noncollinear phase of two-dimensional antiferromagnetic topological insulator MnBi₄Te₇, *Chem. Mater.* **33**, 8343 (2021).
- [21] A. Löhken, C. Lux, D. Johrendt, and A. Mewis, Kristall- und elektronische strukturen von AlR₂P₂ (A: Ca - Ba), *Z. Anorg. Allg. Chem.* **628**, 1472 (2002).
- [22] D. J. García, V. Vildosola, A. A. Aligia, D. G. Franco, and P. S. Cornaglia, Magnetic properties of chiral EuIr₂P₂, *Phys. Rev. B* **104**, 214411 (2021).
- [23] D. G. Franco and C. Geibel, Synthesis and study of the chiral magnetic system EuIr₂P₂, *Phys. Rev. B* **104**, 054416 (2021).
- [24] See Supplemental Material at <http://link.aps.org/supplemental/10.1103/PhysRevResearch.6.043295> for crystal growth optimization, additional single-crystal neutron diffraction, additional magnetic and transport properties, and band structure calculation.
- [25] G. M. Sheldrick and T. R. Schneider, SHELXL: High-resolution refinement, in *Methods in Enzymology* (Elsevier, New York, 1997), Vol. 277, pp. 319–343.
- [26] S. Mühlbauer, B. Binz, F. Jonietz, C. Pfleiderer, A. Rosch, A. Neubauer, R. Georgii, and P. Böni, Skyrmion lattice in a chiral magnet, *Science* **323**, 915 (2009).
- [27] B. I. Shklovskii and A. L. Efros, *Electronic Properties of Doped Semiconductors*, Softcover reprint of the original 1st edition 1984 (Springer, Berlin, 2013).
- [28] N. F. Mott, Conduction in glasses containing transition metal ions, *J. Noncryst. Solids* **1**, 1 (1968).
- [29] V. Ambegaokar, B. I. Halperin, and J. S. Langer, Hopping conductivity in disordered systems, *Phys. Rev. B* **4**, 2612 (1971).
- [30] R. Rosenbaum, Crossover from Mott to Efros-Shklovskii variable-range-hopping conductivity in In_xO_y films, *Phys. Rev. B* **44**, 3599 (1991).
- [31] Z. J. Razi, S. A. Sebt, and A. Khajehnezhad, Magnetoresistance temperature dependence of LSMO and LBMO perovskite manganites, *J. Theor. Appl. Phys.* **12**, 243 (2018).
- [32] J. Y. Chan, S. M. Kauzlarich, P. Klavins, R. N. Shelton, and D. J. Webb, Colossal negative magnetoresistance in an antiferromagnet, *Phys. Rev. B* **57**, R8103(R) (1998).
- [33] S. Noguchi, M4 scaling of negative magnetoresistance in CeSi, *Physica B* **329**, 491 (2003).
- [34] S. Nakatsuji, N. Kiyohara, and T. Higo, Large anomalous Hall effect in a non-collinear antiferromagnet at room temperature, *Nature (London)* **527**, 212 (2015).
- [35] A. K. Nayak *et al.*, Large anomalous Hall effect driven by a nonvanishing Berry curvature in the noncollinear antiferromagnet Mn₃Ge, *Sci. Adv.* **2**, e1501870 (2016).
- [36] V. Kumar, N. Kumar, M. Reehuis, J. Gayles, A. S. Sukhanov, A. Hoser, F. Damay, C. Shekhar, P. Adler, and C. Felser, Detection of antiskyrmions by topological Hall effect in Heusler compounds, *Phys. Rev. B* **101**, 014424 (2020).
- [37] M. Uchida *et al.*, Above-ordering-temperature large anomalous Hall effect in a triangular-lattice magnetic semiconductor, *Sci. Adv.* **7**, eabl5381 (2021).
- [38] Y. Shapira, S. Foner, N. F. Oliveira, and T. B. Reed, EuTe. II. Resistivity and Hall effect, *Phys. Rev. B* **5**, 2647 (1972).
- [39] A. A. Burkov and L. Balents, Anomalous Hall effect in ferromagnetic semiconductors in the hopping transport regime, *Phys. Rev. Lett.* **91**, 057202 (2003).
- [40] W. Allen, E. G. Gwinn, T. C. Kreutz, and A. C. Gossard, Anomalous Hall effect in ferromagnetic semiconductors with hopping transport, *Phys. Rev. B* **70**, 125320 (2004).
- [41] I. Shlimak, M. Kaveh, M. Yosefin, M. Lea, and P. Fozooni, Crossover phenomenon for hopping conduction in strong magnetic fields, *Phys. Rev. Lett.* **68**, 3076 (1992).
- [42] C. Zeng, Y. Yao, Q. Niu, and H. H. Weitering, Linear magnetization dependence of the intrinsic anomalous Hall effect, *Phys. Rev. Lett.* **96**, 037204 (2006).
- [43] Y. Lyanda-Geller, S. H. Chun, M. B. Salamon, P. M. Goldbart, P. D. Han, Y. Tomioka, A. Asamitsu, and Y. Tokura, Charge transport in manganites: Hopping conduction, the anomalous Hall effect, and universal scaling, *Phys. Rev. B* **63**, 184426 (2001).
- [44] T. Holstein, Hall Effect in impurity conduction, *Phys. Rev.* **124**, 1329 (1961).
- [45] J. M. Perez-Mato, S. V. Gallego, E. S. Tasci, L. Elcoro, G. De La Flor, and M. I. Aroyo, Symmetry-based computational tools

- for magnetic crystallography, *Annu. Rev. Mater. Res.* **45**, 217 (2015).
- [46] G. Kresse and J. Furthmüller, Efficient iterative schemes for *ab initio* total-energy calculations using a plane-wave basis set, *Phys. Rev. B* **54**, 11169 (1996).
- [47] G. Kresse and D. Joubert, From ultrasoft pseudopotentials to the projector augmented-wave method, *Phys. Rev. B* **59**, 1758 (1999).
- [48] P. E. Blöchl, Projector augmented-wave method, *Phys. Rev. B* **50**, 17953 (1994).
- [49] J. P. Perdew, K. Burke, and M. Ernzerhof, Generalized gradient approximation made simple, *Phys. Rev. Lett.* **77**, 3865 (1996).
- [50] V. I. Anisimov, F. Aryasetiawan, and A. I. Lichtenstein, First-principles calculations of the electronic structure and spectra of strongly correlated systems: The **LDA**+*U* method, *J. Phys.: Condens. Matter* **9**, 767 (1997).
- [51] K. Ueno, T. Fukumura, H. Toyosaki, M. Nakano, and M. Kawasaki, Anomalous Hall effect in anatase $\text{Ti}_{1-x}\text{Co}_x\text{O}_{2-\delta}$ at low temperature regime, *Appl. Phys. Lett.* **90**, 072103 (2007).
- [52] H. Ohno, H. Munekata, T. Penney, S. Von Molnár, and L. L. Chang, Magnetotransport properties of *p*-type (In,Mn)As diluted magnetic III-V semiconductors, *Phys. Rev. Lett.* **68**, 2664 (1992).
- [53] H. Ohno, A. Shen, F. Matsukura, A. Oiwa, A. Endo, S. Katsumoto, and Y. Iye, (Ga,Mn)As: A new diluted magnetic semiconductor based on GaAs, *Appl. Phys. Lett.* **69**, 363 (1996).
- [54] G. Mihály, M. Csontos, S. Bordács, I. Kézsmárki, T. Wojtowicz, X. Liu, B. Jankó, and J. K. Furdyna, Anomalous Hall effect in the (In,Mn)Sb dilute magnetic semiconductor, *Phys. Rev. Lett.* **100**, 107201 (2008).
- [55] M. Eginligil, G. Kim, Y. Yoon, J. P. Bird, H. Luo, and B. D. McCombe, Manipulation of an unusual anomalous Hall effect in $\text{Ga}_{1-x}\text{Mn}_x\text{Sb}$ random alloys, *Physica E* **40**, 2104 (2008).
- [56] D. Chiba, A. Werpachowska, M. Endo, Y. Nishitani, F. Matsukura, T. Dietl, and H. Ohno, Anomalous Hall effect in field-effect structures of (Ga,Mn)As, *Phys. Rev. Lett.* **104**, 106601 (2010).
- [57] A. Asamitsu and Y. Tokura, Hall effect in $\text{La}_{1-x}\text{Sr}_x\text{MnO}_3$, *Phys. Rev. B* **58**, 47 (1998).
- [58] P. Matl, N. P. Ong, Y. F. Yan, Y. Q. Li, D. Studebaker, T. Baum, and G. Doubinina, Hall effect of the colossal magnetoresistance manganite $\text{La}_{1-x}\text{Ca}_x\text{MnO}_3$, *Phys. Rev. B* **57**, 10248 (1998).
- [59] W.-L. Lee, S. Watauchi, V. L. Miller, R. J. Cava, and N. P. Ong, Dissipationless anomalous Hall current in the ferromagnetic spinel $\text{CuCr}_2\text{Se}_{4-x}\text{Br}_x$, *Science* **303**, 1647 (2004).
- [60] S. U. Yuldashev, H. C. Jeon, H. S. Im, T. W. Kang, S. H. Lee, and J. K. Furdyna, Anomalous Hall effect in insulating $\text{Ga}_{1-x}\text{Mn}_x\text{As}$, *Phys. Rev. B* **70**, 193203 (2004).