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Phys. Rev. Lett. **134**, 186501 — Published 6 May 2025

DOI: [10.1103/PhysRevLett.134.186501](https://doi.org/10.1103/PhysRevLett.134.186501)

Anomalous Electrical Transport in the Kagome Magnet YbFe_6Ge_6

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(Dated: April 15, 2025)

Two-dimensional (2D) kagome metals offer a unique platform for exploring electron correlation phenomena derived from quantum many-body effects. Here, we report a combined study of electrical magnetotransport and neutron scattering on YbFe_6Ge_6 , where the Fe moments in the 2D kagome layers exhibit an *A*-type collinear antiferromagnetic order below $T_N \approx 500$ K. Interactions between the Fe ions in the layers and the localized Yb magnetic ions in between reorient the *c*-axis aligned Fe moments to the kagome plane below $T_{SR} \approx 63$ K. Our magnetotransport measurements show an intriguing anomalous Hall effect (AHE) that emerges in the spin-reorientated collinear state, accompanied by the closing of the spin anisotropy gap as revealed from inelastic neutron scattering. The gapless spin excitations and the Yb-Fe interaction are able to support a dynamic scalar spin chirality, which explains the observed AHE. Therefore, our study demonstrates spin fluctuations may provide an additional scattering channel for the conduction electrons and give rise to AHE even in a collinear antiferromagnet.

The Hall effect, discovered by Edwin Hall, is the potential difference across an electric conductor transverse to the current and to an externally applied magnetic field perpendicular to the current [1]. For conventional metals or semiconductors, the ordinary Hall effect (voltage) is linearly proportional to the applied magnetic field. The anomalous Hall effect (AHE), one of the most prominent phenomena of quantum transport in correlated electron materials, describes unconventional electron deflection other than the effect of external magnetic fields [2–4]. While the initial observation of AHE was made in ferromagnets (conventional AHE) [2], AHE can also occur in materials without net magnetization. The observation of AHE in antiferromagnets is of particular interest, where the interplay of itinerant electrons and spin texture (i.e., magnetic structure) gives rise to this unique transport property [5–7]. From the application perspective, antiferromagnetic (AFM) spintronic devices have advantages over ferromagnetic counterparts, such as the absence of

stray fields and being robust to external magnetic field perturbation [8].

Therefore, a determination of the microscopic origin of the AHE in antiferromagnets has been attracting great interest. For example, it is well-known that a real-space Berry phase originating from a magnetic field-induced skyrmion lattice (non-coplanar spin texture) can produce an AHE termed “topological Hall effect” (THE) [9–14]. The non-coplanar spin texture of a skyrmion has a non-zero scalar spin chirality (SSC) $\chi_{ijk} = \mathbf{S}_i \cdot (\mathbf{S}_j \times \mathbf{S}_k)$, where $\mathbf{S}_i$, $\mathbf{S}_j$, and $\mathbf{S}_k$ are three localized spins at sites *i*, *j*, and *k*, that acts as a fictitious magnetic field to induce AHE [4]. The chiral spin texture of a magnetic skyrmion is often stabilized by the competition of Heisenberg and Dzyaloshinskii-Moriya interactions [9–11].

Going beyond the paradigm of the static spin effect, recent studies suggest that fluctuating spins can also contribute to the AHE by scattering the electrons, even in the spin-disordered state [15–18]. Arranging spins in a triangular motif is considered to be a favorable condition [18], as spin fluctuations could be significantly promoted by the geometrical frustration. The kagome lattice, a two-dimensional (2D) network of corner-sharing triangles, is an ideal platform to realize this mechanism.

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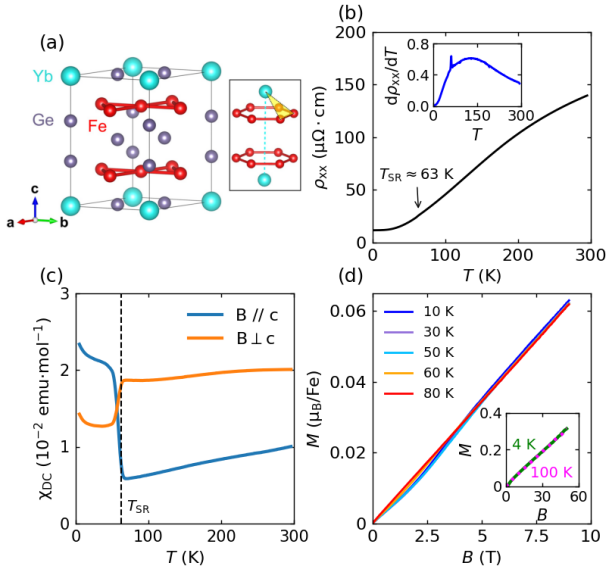


FIG. 1. (a) Crystal structure of YbFe_6Ge_6 . Right inset shows the local structure of Yb and Fe, with one spin triad of Yb-Fe highlighted in yellow. (b) Temperature dependence of zero-field resistivity in the kagome plane. Inset shows its derivative with respect to temperature. (c) Temperature dependence of DC magnetic susceptibility under a magnetic field of 1 T parallel and perpendicular to the c -axis. (d) Isothermal magnetizations along the $\mathbf{a}^*$ direction up to 9 T at selected temperatures. Inset shows the magnetizations up to 50 T at 4 K and 100 K.

Indeed, spin-fluctuation driven AHE/THE was recently observed in some kagome magnets, including AMn_6Sn_6 ($A = \text{Y, Sc, and Er}$) [19–22], $\text{B}_3\text{Ru}_4\text{Al}_{12}$ ($B = \text{Nd and Gd}$) [23, 24], and HoAgGe [25]. In these materials, non-zero time-averaged SSC $\langle \mathbf{S}_i \cdot (\mathbf{S}_j \times \mathbf{S}_k) \rangle$ is induced by spin fluctuations out of noncollinear magnetic structures or paramagnetic states. However, there is no direct evidence of such spin fluctuations, for example, by inelastic neutron scattering (INS) experiments, particularly in the temperature region where the AHE was detected. In addition, it is still unclear whether spin fluctuations from collinear AFM structures can have $\langle \mathbf{S}_i \cdot (\mathbf{S}_j \times \mathbf{S}_k) \rangle \neq 0$ and therefore AHE/THE.

Here we report the discovery of a spin-fluctuation driven AHE in the kagome lattice magnet YbFe_6Ge_6 with a collinear AFM structure [26–30]. Through a comprehensive study of magnetization, magnetotransport, and neutron diffraction on single crystals, we find that the AHE emerges in its spin-reorientated state below $T_{\text{SR}} \approx 63$ K, where the collinear Fe spins rotate from the c -axis to the kagome plane. The combined space inversion and time-reversal ($\mathcal{IT}$) symmetry of the magnetic structure rules out the static non-zero SSC being the origin of the AHE [5–7]. Since our INS experiments reveal gapless low-energy spin excitations in the spin-reorientated state, the highly fluctuating Fe spins can contribute to a dynamic non-zero SSC by interacting with the Yb spins, which thereby induces the AHE [31].

The crystal structure of YbFe_6Ge_6 is shown in Fig. 1(a), where the Fe ions form a 2D kagome lattice. YbFe_6Ge_6 belongs to the family of hexagonal $R\text{Fe}_6\text{Ge}_6$ compounds (with $R = \text{Sc, Yb, Lu, Mg, Ti, Zr, Hf, and Nb}$) [26, 28, 32–34]. Their structure is derived by inserting the R ions into the B35 structure FeGe [35]. Despite the fact that they all have a Néel temperature well above room temperature, YbFe_6Ge_6 is the only member exhibiting a spontaneous spin reorientation (SR) transition [26, 28, 32–34]. Yb has been determined to be trivalent, yet its triangular lattice does not order at least above 0.4 K [27, 28, 36]. Thus the magnetism of YbFe_6Ge_6 is mainly given by the Fe kagome lattice, although there must also be Yb-Fe interactions at low temperatures.

The zero-field resistivity of YbFe_6Ge_6 along the kagome plane is presented in Fig. 1(b), which shows a typical metallic behavior with a residual-resistivity ratio of ~ 12 . The derivative of the resistivity shows a small yet sharp peak around T_{SR} , indicating the change of magnetic structure. Fig. 1(c) shows the DC magnetic susceptibility under a magnetic field of 1 T along and perpendicular to the c -axis. The sudden susceptibility change around 63 K confirms the occurrence of the SR transition. These electrical and magnetic characteristics are consistent with an early report [28].

Next, we explore how external fields affect the magnetic and transport properties of YbFe_6Ge_6 . Fig. 1(d) presents the isothermal magnetization along the $\mathbf{a}^*$ direction at selected temperatures from 10 K to 100 K. Although the magnetization shows an overall linear dependence on the field, it is slightly nonlinear at low fields for $T < T_{\text{SR}}$. However, the magnitude of magnetization is quite small. Even at a magnetic field of 50 T, the moment size is only about $0.3 \mu_{\text{B}}/\text{Fe}$ [inset of Fig. 1(d)], which is much smaller than the full moment size of $\sim 1.5 \mu_{\text{B}}/\text{Fe}$ [26, 27, 29, 30, 36]. The contrast between $T > T_{\text{SR}}$ and $T < T_{\text{SR}}$ can be more clearly discerned in the magnetoresistance (MR) [Fig. 2(a)], with $\text{MR} = [\rho_{xx}(B) - \rho_{xx}(0 \text{ T})]/\rho_{xx}(0 \text{ T})$. Here the magnetic field is also applied along the $\mathbf{a}^*$ direction. At 80 K, the MR has an approximately quadratic dependence on the field up to 12 T. When cooling below T_{SR} , it first increases rapidly at low field and then shows the approximately quadratic behavior at high fields. Similar MR behavior was observed in Fe_3Sn_2 , which is typical of metallic magnets with an SR transition [37, 38].

Field dependence of the Hall resistivity is shown in Fig. 2(b), where we find a drastic change of ρ_{yx} from 10 K to 80 K. In magnetic systems, the Hall resistivity can be decomposed into three parts: $\rho_{yx} = \rho_{yx}^{\text{O}} + \rho_{yx}^{\text{A}} + \Delta\rho_{yx}$ [31, 39], where ρ_{yx}^{O} is ordinary Hall resistivity, ρ_{yx}^{A} is conventional anomalous Hall resistivity due to net magnetization ($\rho_{yx}^{\text{A}} \propto M$), and $\Delta\rho_{yx}$ is additional AHE contributions that cannot be included in the former two parts. For a magnetic system with localized moments, $\Delta\rho_{yx}$ usually comes from static and/or dynamic spin textures. In our case, M is very small and shows barely temperature dependence [Fig. 1(d)], which is in contrast to the

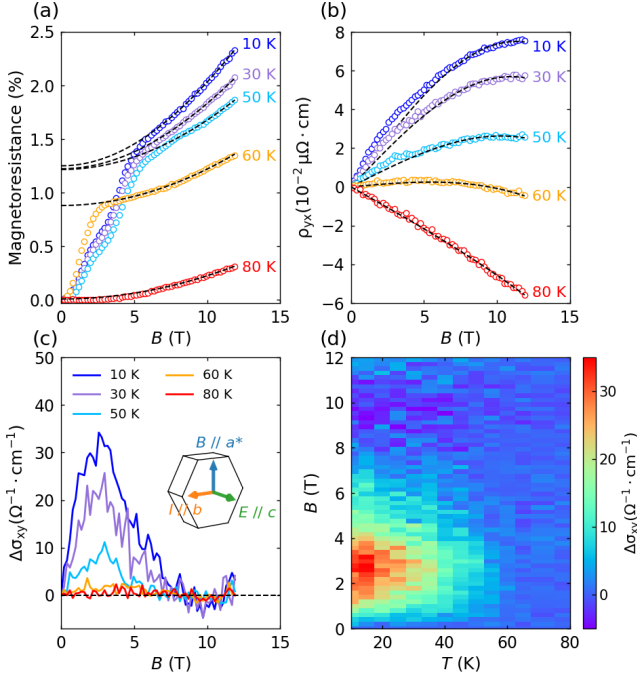


FIG. 2. (a) Magnetoresistance at selected temperatures under fields along the $\mathbf{a}^*$ direction. The electric current is along the b -axis. (b) Hall resistivity at selected temperatures under fields along the $\mathbf{a}^*$ direction. Dashed curves in (a) and (b) are the fits with the two-band model [36]. (c) Remnant Hall conductivity ($\Delta\sigma_{xy}$) after subtracting the ordinary Hall contribution. Inset shows the Hall measurement configuration. (d) Contour plot of $\Delta\sigma_{xy}$ for temperatures from 10 K to 80 K and magnetic fields from 0 T to 12 T.

temperature evolution of ρ_{yx} [Fig. 2(b)]. Therefore, ρ_{yx}^A is negligible, and we focus on the remaining two parts.

The nonlinear field dependence of MR and ρ_{yx} indicates a dominant multiband behavior. We use a two-band model to simultaneously fit the MR and ρ_{yx} [38, 40, 41]. Although the data at high fields can be well described by the two-band model, evident deviation from it at the low field region can be observed below T_{SR} [Fig. 2(a) and (b)], indicating a contribution from $\Delta\rho_{yx}$. Fig. 2(c) shows the corresponding Hall conductivity $\Delta\sigma_{xy}$, where it only appears below T_{SR} . With more detailed measurements at intermediate temperatures [36], we construct a contour plot of $\Delta\sigma_{xy}$, which shows a remarkable similarity to those metallic magnets with novel Hall effects [10, 39]. The maximum magnitude of the Hall conductivity at 10 K is about $30 \Omega^{-1}\text{cm}^{-1}$, comparable with those in previous “anomalous Hall antiferromagnets” [7]. For magnetic fields applied perpendicular to the kagome plane, however, we do not observe such an AHE [36], suggesting its intimate connection to the magnetic anisotropy as will be discussed below.

As the anomalous electrical transport is closely related to the SR, knowing the change of the magnetic structure across T_{SR} is crucial to uncovering its origin. We hence use single crystal neutron diffraction to study

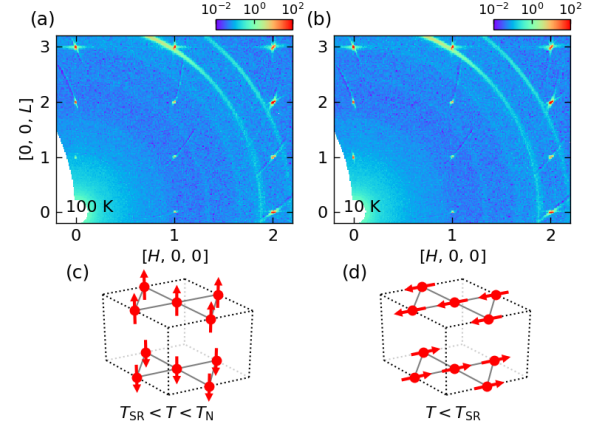


FIG. 3. (a) and (b) Neutron diffraction patterns of the $(H, 0, L)$ plane at 100 K and 10 K, respectively. The intensity color bar is in log scale to highlight possible weak features. (c) and (d) Magnetic structures of the Fe kagome lattice above and below T_{SR} , respectively.

the magnetic structure of YbFe_6Ge_6 . Fig. 3(a) shows the diffraction pattern of the $(H, 0, L)$ plane at 100 K ($> T_{SR}$). There are no additional Bragg peaks except those at integer-index positions, consistent with the well-established A-type AFM structure with a propagation wave vector $\mathbf{k}_m = (0, 0, 0)$ [Fig. 3(c)] [27, 29, 30]. When cooling down to 10 K ($< T_{SR}$), the diffraction pattern is strikingly similar [Fig. 3(b)] - no additional peaks appear, while the peak intensities have clear temperature-dependent behavior due to the SR [36]. This suggests the $\mathbf{k}_m$ of YbFe_6Ge_6 remains to be $(0, 0, 0)$ below T_{SR} . Representation analysis based on this propagation wave vector shows only a collinear structure with the spins along the a -axis [Fig. 3(d)] is compatible with the experiment. The SR transition in YbFe_6Ge_6 therefore equivalently flips the Néel vector by 90° . More details of the magnetic structure analysis can be found in [36]. The magnetic structure below T_{SR} along with the lattice space group $P6/mmm$ ensures the system has the $\mathcal{IT}$ symmetry. With small-angle neutron scattering, we also confirmed no field-induced peaks appear under magnetic fields along the $\mathbf{a}^*$ direction [36]. These observations rule out that the observed AHE is induced by static spin textures, such as magnetic skyrmions [13], which are restricted by Berry curvature that requires the absence of $\mathcal{IT}$ symmetry [4, 6, 42]. Conversely, the spin-fluctuation driven AHE (or THE) is not necessarily dictated by this symmetry, as it originates from electron scattering by dynamic spins [15–18].

Since YbFe_6Ge_6 has localized magnetic moments, spin fluctuations below T_N are spin waves (magnons), which could impact the motion of conduction electrons [15]. Hence, we use INS to directly probe the magnons and determine their temperature evolution. Fig. 4(a) shows the energy dependence of the INS intensities at $(0, 0, 1)$. At 2.5 K, some INS intensities extend to ~ 0.6 meV, which is much larger than our instrumental energy resolution

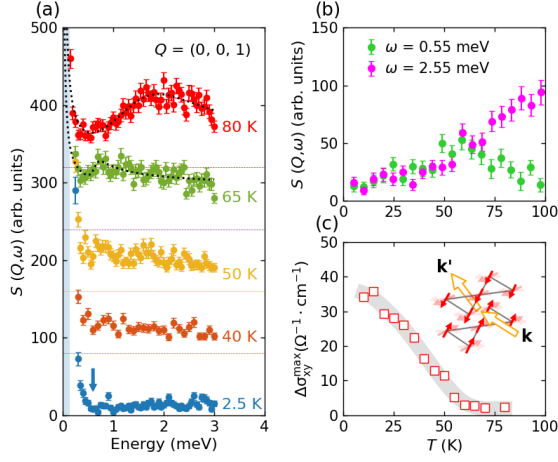


FIG. 4. (a) Low-energy spin excitations at (0, 0, 1) at selected temperatures, offset for clarity. Horizontal dashed lines indicate the zero intensity for the data above 2.5 K. Dashed curves are the fits to 65 K and 80 K data [36]. The light blue region shows energy resolution, and the arrow marks 0.6 meV, below which gapless excitations emerge. (b) Temperature dependence of the intensities at (0, 0, 1) at 0.55 meV and 2.55 meV. (c) Temperature dependence of the maximum magnitudes of $\Delta\sigma_{xy}$. The bold grey curve is a guide to the eyes. Inset illustrates electron scattering by the spin fluctuations.

(~ 0.14 meV). Therefore, the spin excitations are gapless in nature. As the temperature increases, their intensity significantly increases, but they remain gapless until 50 K. When the temperature goes beyond T_{SR} , the intensity of the low-energy part starts to decrease. A broad peak centering around 2 meV shows up at 80 K, indicating the opening of a spin-wave gap. Such gap opening behavior suggests the magnetic anisotropy of YbFe_6Ge_6 transfers from an easy-plane type below T_{SR} to an easy-axis type above T_{SR} , consistent with the magnetic structures presented in Fig. 3(c) and (d). By fitting the data with an error function multiplied by the Bose factor [36, 43], we can extract gap sizes of 0.64(4) meV and 1.34(5) meV at 65 K and 80 K, respectively.

To determine the evolution of the spin-wave gap, we show the temperature dependence of the intensities at (0, 0, 1) at 0.55 meV and 2.55 meV [Fig. 4(b)], which are located below and above the spin-wave gap, respectively. Below T_{SR} , the intensities at both energies increase as the temperature goes up. However, they start to bifurcate at T_{SR} - while the intensity above the spin-wave gap continues increasing, the intensity below it starts to decrease. The continuous increase of the intensity at 0.55 meV below T_{SR} further suggests it is an excitation signal rather than the tail of Bragg peak, whose intensity should otherwise decrease when the temperature approaches T_{SR} . The onset of gapless spin excitations below T_{SR} aligns with the AHE [see Fig. 4(c)], suggesting their correlation.

We point out that the gapless spin excitations below T_{SR} play a vital role in explaining the AHE through a dy-

namic SSC scattering mechanism, as first identified in the ferromagnetic heterostructure $\text{SrRuO}_3/\text{SrTiO}_3$ [31]. The gapless excitations at (0, 0, 1) within our energy resolution suggests that Fe spins can fluctuate at the cost of a very small energy, facilitating the formation of a dynamic SSC. From a semiclassical perspective, the propagation of spin waves causes adjacent Fe spins in the kagome plane to slightly deviate from the collinear ground state. On the other hand, Yb spins remain disordered above 0.4 K and thus, by themselves, have little impact on the AHE, which emerges at a much higher temperature. However, their interactions with Fe strengthen upon cooling, leading to the SR transition. Because the Yb positions and Fe kagome layers are well separated [Fig. 1(a)], the transient non-collinear Fe spin configurations and Yb spins can locally produce an SSC with $\mathbf{S}_i \cdot (\mathbf{S}_j \times \mathbf{S}_k) \neq 0$. Despite these local SSCs with opposite signs (+ or -) cancel out at zero magnetic field due to the global $\mathcal{IT}$ symmetry, a finite magnetic field partially aligns the Yb spins, resulting in a non-zero overall SSC [36]. This non-zero SSC can preferentially scatter electrons and give rise to the observed AHE. However, when the magnetic field becomes exceedingly large, low-energy spin fluctuations are suppressed due to the opening of a Zeeman gap, leading to a vanishing overall SSC and, consequently, a zero AHE. In contrast, the finite spin-wave gap above T_{SR} suggests relatively high spin stiffness. As a result, Fe spins favor maintaining the ground-state collinear configuration. Both the increased spin stiffness of the Fe kagome layer and the weaker Yb-Fe interaction at higher temperatures are unfavorable for the dynamic SSC, which accounts for the absence of the AHE above T_{SR} .

Therefore, the AHE in YbFe_6Ge_6 can be understood as arising from electron scattering by low-energy spin fluctuations. Given that every Fe ion has a static magnetic moment of approximately $1.5 \mu_B$ below T_{SR} [36], we estimate a Zeeman gap of roughly 0.6 meV for a 7 T field [43, 44], beyond which the AHE vanishes at low temperatures [Fig. 2(c) and (d)]. It suggests the dynamic SSC contributing to the AHE is within 0.6 meV. In addition, this energy scale reasonably matches the gap size at 65 K [Fig. 4(a)], which is slightly above T_{SR} and the AHE just disappears [Fig. 4(c)]. For comparison, we note that the similar kagome antiferromagnet FeSn does not exhibit an AHE [45, 46]. Despite also having a collinear magnetic structure with spins lying in the kagome plane, the spin waves of FeSn display a finite energy gap of ~ 2 meV [47, 48], which suppresses low-energy spin fluctuations.

Recent experiments show that spin fluctuations can give rise to similar AHEs (or THEs) in a range of materials, including single crystals [19–25, 49, 50] and thin films [31, 51], supporting the idea that spin-fluctuation driven AHE is a general phenomenon. In previous studies, it depends on thermal fluctuations, with samples being heated above or close to the ordering temperature [19–25, 31, 49–51]. Those AHEs are therefore essentially driven by paramagnetic spin fluctuations, which are usually dominant around magnetic Bragg peak positions near the transition

temperature and are gapless as well [52]. For YbFe_6Ge_6 , the situation is different. Due to the SR transition induced by the Yb-Fe interaction, the c -axis aligned spins transfer to the easy-plane arrangement below T_{SR} , which greatly promotes low-energy spin fluctuations even in the spin-ordered state. With the stronger Yb-Fe interaction at low temperatures, the AHE is enhanced accordingly. Moreover, YbFe_6Ge_6 adopts a collinear AFM structure that is much simpler than the magnetic structures of previous single crystals [19–25, 49, 50]. Our study therefore also suggests that a complex (e.g., chiral) static magnetic structure is not a prerequisite for observing spin-fluctuation driven AHE.

In conclusion, with electrical transport, magnetization, and neutron scattering, we have systematically studied the Fe-based kagome magnet YbFe_6Ge_6 . An AHE is observed in its AFM-ordered state when the spins are aligned parallel to the kagome plane due to the Yb-Fe interaction. The collinear magnetic structure with $\mathcal{IT}$ symmetry excludes a static SSC origin of the AHE. By directly measuring the low-energy spin excitations, we find that spin excitations at the Brillouin zone center become gapless in the spin-reorientated state. The simultaneous onset of these gapless excitations and AHE shows their close relationship and points to a dynamic SSC scattering mechanism of the latter. Our results not only reveal a spin-fluctuation driven AHE that is not solely dependent on thermal activation but also demonstrate dynamic spins, even in a collinear antiferromagnet,

can significantly impact electrical transport.

ACKNOWLEDGMENTS

The neutron scattering and single-crystal synthesis work at Rice was supported by US NSF-DMR-2401084 and by the Robert A. Welch Foundation under grant no. C-1839 (P.D.), respectively. The transport work at University of Tokyo was supported by Grant-in-Aid for Scientific Research (KAKENHI) (JP22H00105), and Grant-in-Aid for Scientific Research on Innovative Areas “Quantum Liquid Crystals” (JP19H05824) from Japan Society for the Promotion of Science. W.Y. acknowledges support from the US NSF-OISE-2201516 under the Accelerated program of Office of International Science and Engineering (OISE). P.D. also acknowledges travel support from the US NSF-OISE-2201516. A portion of this research used resources at the Spallation Neutron Source, a DOE Office of Science User Facility operated by the Oak Ridge National Laboratory. The experiment at the ISIS Neutron and Muon Source was supported by a beam time allocation RB2410134 from the Science and Technology Facilities Council. Ø.S.F. acknowledges funding from The Research Council of Norway through project 325345. This work is in part based on work performed at the DMC instrument at SINQ at the Paul Scherrer Institute. Small-angle neutron scattering was performed at QUOKKA, ANSTO under a proposal number P18792. CLH acknowledges the support by the National Science and Technology Council in Taiwan (NSTC 113-2112-M-006-027).

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End Matter

Appendix – While dispersive spin waves of YbFe_6Ge_6 extend to much higher energies, only the low-energy part (i.e., the spin-wave gap) shows a significant change across T_{SR} . Fig. 5(a) and (b) show the excitation spectra along the $[0, 0, L]$ direction at 10 K and 100 K, respectively [36]. The signals around $L = 1$ and 3 are magnons, while phonons appear at larger momentum transfer positions. The overall excitation spectrum changes little except for the intensity, as can be seen from the constant-energy cuts shown in Fig. 5(c) and (d). We notice the intensity increase from 10 K to 100 K is significantly larger than the effect simply due to the Bose factor [grey curves in Fig. 5(c) and (d)], indicating a change of the spin orientation. Thermal evolution of the low-energy spin excitations around $(0, 0, 1)$ is presented in Fig. 5(e)-(i) [36]. Since the magnon has a bandwidth of ~ 40 meV along the $[0, 0, L]$ direction [36], the spin excitations below 3 meV are rod-like and concentrated at the Brillouin zone center. As the temperature rises to near T_{SR} , their intensity gradually accumulates, which extends down to the elastic background [Fig. 5(e)-(h)]. However, the intensity around 1 meV drops at 80 K ($> T_{\text{SR}}$) [Fig. 5(i)], indicating the presence of an energy gap, consistent with

the data shown in Fig. 4.

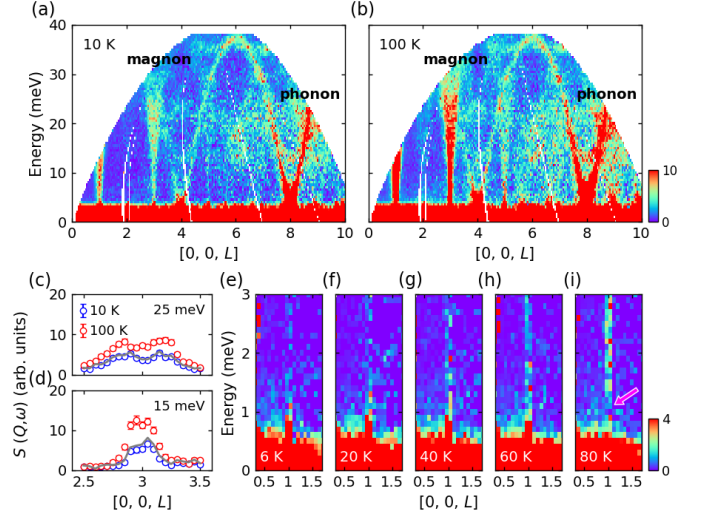


FIG. 5. (a) and (b) Excitation spectra along $[0, 0, L]$ direction at 10 K and 100 K, respectively. (c) and (d) Constant energy cuts around $(0, 0, 3)$ at 25 meV and 15 meV, respectively. Grey curves represent the calculated intensities at 100 K by correcting the 10 K data with the Bose factor. (e)-(i) Low-energy spin excitations around $(0, 0, 1)$ at selected temperatures. Magenta arrow indicates the energy gap at 80 K.