

Topological surface magnon-polariton in an insulating canted antiferromagnet

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Excitation and control of antiferromagnetic magnon modes lie at the heart of coherent antiferromagnetic spintronics. Here, we propose a topological surface magnon-polariton as a new approach in the prototypical magnonic material hematite. We show that in an insulating canted antiferromagnet, where strong-coupled magnon-photon modes can be achieved using on-chip microwave layouts, a surface magnon-polariton mode exists in the gap of the bulk magnon-photon bands. The emergence of surface magnon-polariton mode is further attributed to the nontrivial topology of bulk magnon-photon bands. Magnon-photon coupling enhances the Berry curvature near the anticrossing points, leading to a topological bulk Chern band associated with the surface magnon-polaritons. Our work provides a general principle for the utilization of optomagnetic properties in antiferromagnets, with an illustration of its experimental feasibility and wide generality as manifested in hematite.

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I. INTRODUCTION

Magnons, or spin waves, are collective excitations of magnetic moments that can carry and propagate spin-precessional information in magnetic materials without charge transport, and therefore can be utilized in spintronic devices without Ohmic loss [1–6]. The application of magnons in electronic devices has given rise to the field of magnonics, which has been intensively investigated based on ferro- and ferrimagnetic materials, such as metals permalloy [7–9], CoFeB [10–12], semimetal $\text{Co}_2\text{MnGa}_x\text{Ge}_{1-x}$ [13], semiconductor CrI_3 [14,15], insulator yttrium iron garnet [16–18], and their heterostructures [19–21]. By contrast, given their (sub-)terahertz eigenfrequencies [5,22–25], antiferromagnetic magnons were mostly inaccessible to electronic control and detection, and significant progress in antiferromagnetic magnonics has been made only in recent years [23–42].

In van der Waals antiferromagnets, gigahertz magnons have been observed [26–29] due to their weak interlayer exchange interactions, despite their relatively high damping. Recently, the rediscovery of the gigahertz antiferromagnetic magnon bands in an insulating canted antiferromagnet, hematite, has brought about a new area of antiferromagnetic magnonics, which possesses the advantages of low damping, high group velocity, and insensitivity to external magnetic disturbances [30–34]. The interplay between strong antiferromagnetic exchange interaction, easy-plane anisotropy, and Dzyaloshinskii-Moriya interaction (DMI) results in coherent magnon modes at the gigahertz range, thus making the electrical detection of magnons in hematite feasible [30–38]. These features, along with an outstandingly low damping coefficient that can be characterized by multiple experimental means [33,38], establish hematite as an ideal platform for antiferromagnetic magnonics. Hence, methods to excite magnon modes in hematite at room temperature are desirable [5,32–34], and we propose here an electrical approach targeting its microwave-photon-induced surface mode.

Motivated by a recent cavity-free electrical observation of propagating magnon-photon strong coupling mode in insulating canted antiferromagnetic hematite [43], we take this material as the platform for our analysis. The observed on-chip magnon-photon coupling in a slab geometry enables the investigation of surface magnon-polariton modes with finite wave vector, in contrast to conventional cavity

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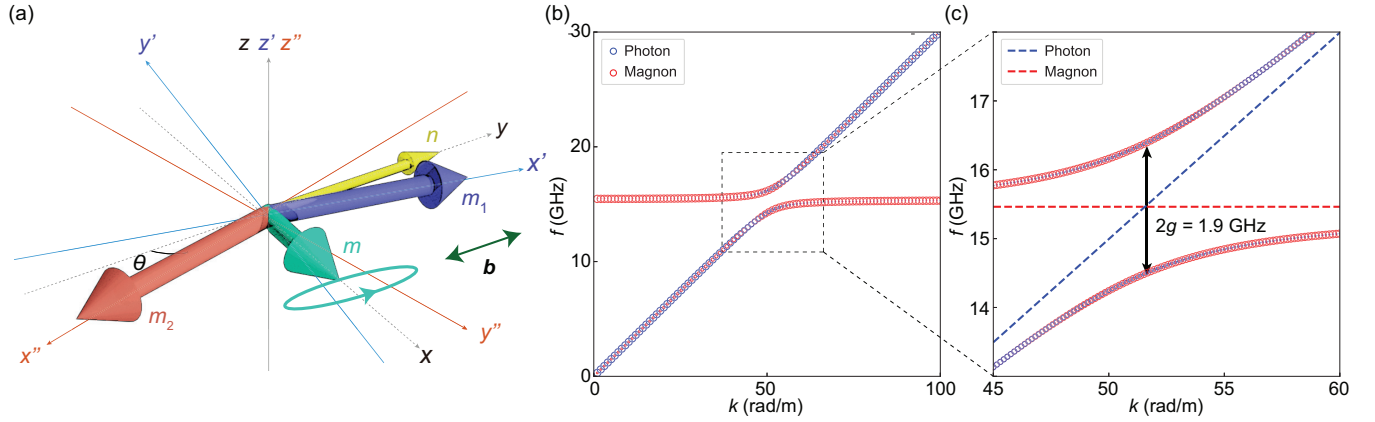


FIG. 1. Modeling bulk magnon-photon bands in the DMI-canted antiferromagnets. (a) The DMI canted model of hematite with the frame of reference setting, sublattices $\mathbf{m}_1, \mathbf{m}_2$, net magnetic moment $\mathbf{m} = (\mathbf{m}_1 + \mathbf{m}_2)/2$, Néel vector $\mathbf{n} = (\mathbf{m}_1 - \mathbf{m}_2)/2$, and an oscillating magnetic field $\mathbf{b}$ coupling to the y -polarized precession. (b) Component-projected dispersion relations of magnon-polaritons. The blue (red) circles denote the photon (magnon) component at the given wave vector and frequency, where the radius of the circle is proportional to the ratio of the corresponding component. (c) Zoomed-in view of the dispersion relations in the region marked by the dashed box in (b), where the blue (red) dashed lines denote the photon (magnon) dispersion without anticrossing. We use $H_0 = 60$ mT, $H_{\text{ex}} = 1040$ T, $H_{\text{DM}} = 2.7$ T, $H_a = 0.067$ mT, $H_A = 1$ mT, and $M_s = 1000$ A/m that give a coupling strength $2g = 1.9$ GHz to qualitatively fit the previous experimental observations [30–32], and in (b),(c), we set $\mathbf{k} \parallel \hat{x}$.

magnonics, which primarily probes bulk resonance modes [44–49]. In easy-axis collinear antiferromagnets, surface magnon-polaritons have been predicted to exist within the magnon-photon coupling gap in the terahertz regime [50]. However, a general physical picture of such an antiferromagnetic surface mode and its correspondence to bulk modes has not yet been established, and it is further addressed here that nontrivial band topology gives rise to the emergence of surface magnon-polaritons in insulating DMI-canted antiferromagnets. Our claim is consistent with the notion of topological magnon-photon anticrossing in ferromagnets [51] and topological magnon-magnon coupling in artificial magnetic lattices [52–55].

In this work, we theoretically investigate the topological surface magnon-polaritons in the insulating, DMI-canted, antiferromagnet hematite. We focus on two aspects: first, whether the surface magnon-polaritons also exist in an easy-plane, DMI-canted antiferromagnetic ground state; second, whether the emergence of antiferromagnetic surface magnon-polaritons is related to the nontrivial topology of the magnon-photon bands. On top of our theoretical discussion, we also provide an experimental protocol for the design of microwave antenna to excite such a surface mode. We use a semiclassical Landau-Lifshitz model combined

with Maxwell’s equations to account for the physics of antiferromagnetic magnons and photons in DMI-canted antiferromagnets. We show the manifestation of this model using hematite as a showcase system, but our model can be extended to any insulating antiferromagnet. The paper is organized as follows: in Sec. II we present a semiclassical model for bulk magnon-photon coupling; in Sec. III we discuss the topology associated with the bulk magnon-photon band gap; in Sec. IV we characterize the surface magnon-polariton modes; and finally in Sec. V we summarize the results and propose experimental protocols.

II. SEMICLASSICAL MODEL FOR BULK MAGNON-PHOTON COUPLING

We first present the semiclassical model for magnon-polaritons in insulating DMI-canted antiferromagnets. In a magnetic media, the modes of electromagnetic waves are governed by the magnetic susceptibility $\chi(\omega)$, which characterizes the motion of the net magnetic moment $\mathbf{m}(\omega)$ under an external excitation magnetic field $\mathbf{b}(\omega)$. To account for the motion of $\mathbf{m}$, we consider a macrospin model with two sublattices $\mathbf{m}_1$ and $\mathbf{m}_2$, which are antiferromagnetically coupled, using the Landau-Lifshitz equation reads (see details in Appendix A) [30–32]

$$\frac{d\mathbf{m}_1}{dt} = -\gamma\mu_0\mathbf{m}_1 \times \left[\mathbf{H}_0 + \frac{\mathbf{b}}{\mu_0} - H_{\text{ex}}\mathbf{m}_2 - H_A(\mathbf{m}_1 \cdot \hat{z})\hat{z} + H_a(\mathbf{m}_1 \cdot \hat{y})\hat{y} + H_{\text{DM}}(\mathbf{m}_2 \times \hat{z}) \right], \quad (1a)$$

$$\frac{d\mathbf{m}_2}{dt} = -\gamma\mu_0\mathbf{m}_2 \times \left[\mathbf{H}_0 + \frac{\mathbf{b}}{\mu_0} - H_{\text{ex}}\mathbf{m}_1 - H_A(\mathbf{m}_2 \cdot \hat{z})\hat{z} + H_a(\mathbf{m}_2 \cdot \hat{y})\hat{y} - H_{\text{DM}}(\mathbf{m}_1 \times \hat{z}) \right], \quad (1b)$$

where $\mathbf{H}_0$ is the applied external magnetic field, $\mathbf{b}/\mu_0$ is the oscillating magnetic field of a photon with μ_0 the vacuum permeability, H_{ex} is the exchange coupling strength between

two sublattices, $H_{A(a)} > 0$ is the out-of-plane (in-plane) hard-axis (easy-plane) anisotropy coefficient whose orientation is $\hat{z}(\hat{y})$, and H_{DM} is the DMI coefficient with the DMI vector

oriented along $\hat{z}$. The frames of reference are illustrated in Fig. 1(a), where the normalized magnetizations are defined as $\mathbf{m}_{1(2)} = \hat{x}^{(n)} + m_{1(2)}^y \hat{y}^{(n)} + m_{1(2)}^z \hat{z}^{(n)}$. We also denote the canting angle induced by DMI as θ for convenience. The exchange stiffness term proportional to $a_{\text{ex}}^2 k^2$ is ignored in this model based on the fact that the wave vector we consider here is much smaller than the reciprocal-lattice constant, i.e., $a_{\text{ex}} k \ll 1$ [32–34]. To capture the eigenmodes of the net magnetic moment, we introduce the transformations $\mathbf{m} = (\mathbf{m}_1 + \mathbf{m}_2)/2$ and $\mathbf{n} = (\mathbf{m}_1 - \mathbf{m}_2)/2$, with $\mathbf{m} = \sin \theta \hat{x} + m_y \hat{y} + m_z \hat{z}$ and $\mathbf{n} = n_x \hat{x} + \hat{y} + n_z \hat{z}$. Accordingly, Eq. (1) is rewritten in terms of $\mathbf{m}$ and $\mathbf{n}$ (see Appendix A), which is feasible for our analysis since the magnon-photon coupling is induced by the Zeeman term $F_{\text{Zeeman}} = -M_s \mathbf{m} \cdot \mathbf{b}$. The full mathematical procedure from magnetic free energy to Landau-Lifshitz equations is further exhibited in Appendix A, which also includes Ref. [32].

Simultaneously, we have the Maxwell's equations accounting for the physics of photons, which read [56]

$$\frac{\partial \mathbf{b}}{\partial t} = -\nabla \times \mathbf{e}, \quad (2a)$$

$$\frac{\partial \mathbf{e}}{\partial t} = \frac{1}{\epsilon} \nabla \times \left(\frac{\mathbf{b}}{\mu_0} - M_s \mathbf{m} \right), \quad (2b)$$

where $\mathbf{e}$ is the electric field of photon, M_s is the magnetization of the canted moment, and ϵ is the dielectric constant taken from experimental results [43]. Here we only consider the displacement current contribution ($\epsilon \partial \mathbf{e} / \partial t$) of electric field in insulators to exclude the irrelevant surface plasmon-polariton solutions [56]. After some algebra (see Appendix A), the effective Schrödinger equation from Landau-Lifshitz and Maxwell's (LLM) equations reads

$$\mathcal{H}_{\text{eff}} \mathbf{x}_k^{(i)} = \omega_i \mathbf{x}_k^{(i)}, \quad (3)$$

where $\mathcal{H}_{\text{eff}}$ is the effective Hamiltonian of bulk magnon-photon modes (see Appendix A), ω_i is the eigenfrequency of the band with index i , and $\mathbf{x}_k^{(i)} \equiv [m_{k,y}^{(i)}, m_{k,z}^{(i)}, n_{k,x}^{(i)}, \mathbf{b}_k^{(i)}, \mathbf{e}_k^{(i)}]^T$ denotes the eigenvector of the band with index i . Here, we only consider the solution with positive frequency, and no imaginary-frequency solution is found. Figures 1(b) and 1(c) show the dispersion relation of lower and upper bulk magnon-photon bands [$f_{1,2} = \omega_{1,2}/(2\pi)$], with parameters given in the caption. We reproduce the observed strong magnon-photon coupling ($2g \sim 2$ GHz) in hematite by adopting a small M_s value, which is also qualitatively compatible with experimental estimates [32].

To normalize the eigenvectors, we define the norm of our eigenvectors as $\|\mathbf{x}\|^2 \equiv \mathbf{x}_k^{(i)\dagger} \Gamma \mathbf{x}_k^{(i)} = \mathcal{E}_i$, where $\mathcal{E}_i$ is the total energy density of the magnon-photon mode with index i , and $\Gamma = \text{diag}[\frac{M_s \omega_m}{\gamma \mu_0} I_{3 \times 3}, \frac{1}{2\mu_0} I_{3 \times 3}, \frac{\epsilon}{2} I_{3 \times 3}]$, with I the unit matrix and $\omega_m = \gamma \mu_0 \sqrt{H_0(H_{\text{DM}} + H_0) + 2H_a H_{\text{ex}}}$ the antiferromagnetic resonance frequency [30–32], following the same strategy as in Ref. [51]. In our calculations, we henceforth choose the normalization condition as $\mathbf{x}_k^{(i)\dagger} \Gamma \mathbf{x}_k^{(i)} = 1$. In Fig. 1(b), we show the component-projected magnon-polariton band structure, where the red (blue) circles indicate the relative contributions of the magnon (photon) components in the wave function by $r = \mathcal{E}_i^{\text{mag(ph)}} / \mathcal{E}_i$, where $\mathcal{E}_i^{\text{mag(ph)}}$ is the magnon (photon) energy density with band index i . The conversion between magnon

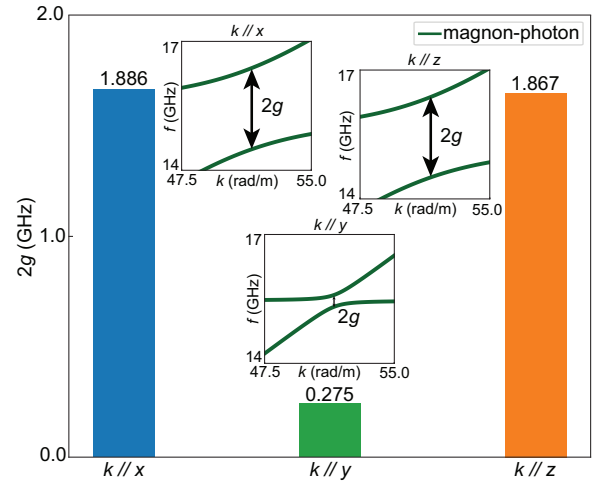


FIG. 2. Anisotropic coupling strength of the bulk magnon-photon bands. Coupling strength $2g$ extracted from bulk magnon-photon band dispersion relations with wave vector along different directions with $\mathbf{k} // \hat{x}$ noted by blue, $\mathbf{k} // \hat{y}$ noted by green, and $\mathbf{k} // \hat{z}$ noted by orange columns, and, respectively, their zoomed-in dispersion relations around the anticrossing point.

and photon components is analogous to the behavior of the hybridized electronic orbitals near the band-inversion-induced anticrossing gap in topological insulators [57–59], where the gap originates from the spin-orbit coupling of electrons. For magnonic anticrossing modes, the coupling associated with its magnetic dipoles with chiral precession can also be treated as a generalized spin-orbit coupling, which further provides nontrivial topology [60–62]. In our model, the eigenmode of the net magnetization $\mathbf{m}$ is a right-handed chiral mode with amplitudes $|m_y| : |m_z| \sim 6:1$ [30,31], as illustrated in Fig. 1(a).

The polarization of $\mathbf{m}$ also leads to an anisotropic magnon-photon coupling strength, as illustrated in Fig. 2. Given that the magnon-photon coupling is mediated by Zeeman coupling ($\propto \mathbf{m} \cdot \mathbf{b}$), its strength is inherently sensitive to whether the polarization direction of $\mathbf{b}$ is aligned with that of $\mathbf{m}$. For example, when the photon wave vector is along the y direction, the coupling interaction is minimized, since the oscillating photon magnetic field aligns in the x - z plane, while the precession of $\mathbf{m}$ is mainly along the y axis. On the contrary, the coupling strength when $\mathbf{k}$ is along x and z is maximized, since the photon magnetic field possesses a y -direction component that couples to $\mathbf{m}$. Therefore, as shown in Fig. 2, the band gap is minimal when the wave vector is along the y direction, whereas for the x - or z -direction cases, the band gaps are approximately five times larger.

Our semiclassical model characterizes the strong magnon-photon coupling in hematite [30,43] even with a small saturated magnetization M_s . If one takes the parameters of other materials with appropriate approximations accordingly, the methodology here can be instantly applied to account for magnon-photon coupling in other insulating antiferromagnetic materials. However, given that each antiferromagnet possesses different eigen precession modes, the involved components in the eigenvector $\mathbf{x}_k^{(i)}$ and the frame-of-reference

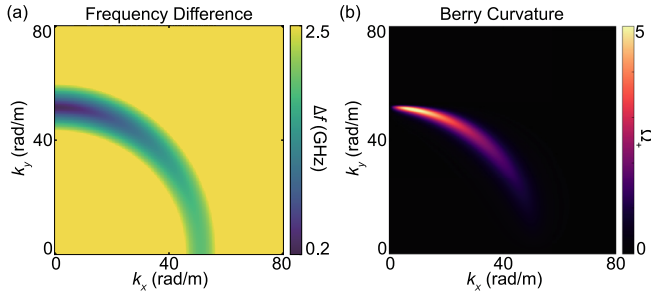


FIG. 3. Distribution of frequency difference ($f_2 - f_1$) and Berry curvature in reciprocal space. (a) Distribution of frequency difference in reciprocal space with a color scale. The minima occur around the anticrossing points. (b) Distribution of the Berry curvature in reciprocal space with $\hbar = 1$. The Berry curvature is enhanced around the minima of bulk band gaps.

transformation should also be adjusted accordingly when applying such an approach to other materials.

III. NONTRIVIAL TOPOLOGY OF BULK MAGNON-PHOTON BANDS

We then discuss the topological properties of the bulk magnon-photon bands. The magnetic ground state of hematite at room temperature is also classified as a weak ferromagnet, characterized by a small but finite net magnetization induced by the DMI [63]. As such, the breaking of time-reversal symmetry inherently occurs in our model. The topological properties of the bulk magnon-phonon bands can be characterized by their Chern numbers [51,52]. First, we propose that the lower bulk magnon-photon band possesses trivial topology. The lower bulk magnon-photon band touches the vacuum photon mode at the Γ point (with zero frequency), which indicates that the lower bulk magnon-photon band is topologically equivalent to the vacuum, or in other words, the lower bulk magnon-photon band carries a zero Chern number $C = 0$. We then investigate the topology of the upper bulk magnon-photon band, which does not touch any other bands in the energy range of our interest. This band could also touch other higher-energy bands of magnons in hematite [64], however the next magnonic band lies around 150 GHz, that is, far beyond the frequency range of the present study, which is centered around 15 GHz [30,31,64]. We thus assume that the topological property of bulk magnon-photon bands only associates with the lowest two magnon-photon bands. We calculate the Berry curvature, which is defined as follows [51,65–69]:

$$\Omega(\mathbf{k}) = -2 \operatorname{Im} \frac{\langle \mathbf{x}_k^{(1)} | \partial_{k_x} \mathcal{H}_{\text{eff}} | \mathbf{x}_k^{(2)} \rangle \langle \mathbf{x}_k^{(2)} | \partial_{k_y} \mathcal{H}_{\text{eff}} | \mathbf{x}_k^{(1)} \rangle}{\hbar(\omega_1 - \omega_2)^2}. \quad (4)$$

The format of Berry curvature is chosen to eliminate the gauge problem [51,66–69]. Given that our focus is on the topological surface magnon-polariton mode at the surface perpendicular to $\hat{z}$, in Eq. (3), and the following analysis, we set $k_z = 0$ in the Hamiltonian. In Fig. 3(a), we show the distribution of frequency difference $\Delta f = f_2 - f_1$ between upper and lower bulk magnon-photon between the upper and lower bulk

magnon-photon band, and in Fig. 3(b) we show the Berry curvature of the upper band. As shown in Fig. 3(a), the frequency difference varies as a function of the wave-vector direction in agreement with our discussion on anisotropic coupling, reaching minima at the anticrossing points. In Fig. 3(b), we can find that for the upper bulk magnon-photon band, the magnitude of Berry curvature Ω_+ is significantly enhanced where the band gap is minimal. This can be understood by the fact that at the anticrossing point with the smallest band gap, the eigenvector switched from magnon-dominated to photon-dominated most rapidly, as demonstrated in Figs. 1(b) and 1(c). The rapid mode shifting provides a large differentiation of the eigenvector $[\partial_{\mathbf{k}} | \mathbf{x}^{(i)} \rangle]$, and further contributes to the local Berry curvature. Also, the shifting of the eigenvector has the potential to provide nontrivial winding in the reciprocal space, as characterized by the Chern number.

We further calculate the Chern number of the upper bulk magnon-photon band by an integration of Berry curvature in the first BZ [51,52]: $C_+ = \int_{\text{BZ}} \Omega(\mathbf{k}) d\mathbf{k} / (2\pi) = 1$. We assume that the boundary of the first BZ is far away from the dispersion we investigated here, and we only integrate around $\mathbf{k}$ points with nonzero Berry curvature [70]. The nonzero Chern number confirms that the upper bulk magnon-photon band is topologically nontrivial. Anticrossing between the bulk magnon and photon modes acts as a source of the Berry curvature, and further gives rise to the nontrivial topology of the upper band, which is the main claim of the nontrivial band topology in this section.

Given the nontrivial upper magnon-photon bulk band alongside with the topologically trivial lower band, we could predict that in the gap of a bulk magnon-photon band there should exist in an emergent mode. In real space, the vacuum is topologically trivial, and the bulk of hematite possesses a topologically nontrivial magnon-photon mode. Thus, at the interface of hematite and vacuum, the aforementioned emergent mode should appear, which can be concluded as a surface mode given our slab geometry. In our system, such discussions on bulk-boundary correspondence can be interpreted as a sign of the existence of a topological surface magnon-polariton, as further discussed in the following text.

IV. SURFACE MAGNON-POLARITON MODE

In this section, we exhibit the surface magnon-polariton mode from our semiclassical approach. Given our prediction, we conduct calculations for the surface magnon-polariton mode from our semiclassical approach. Our methodology follows Ref. [50], in which the authors also found surface magnon-polariton modes in collinear antiferromagnets. We provide the details of calculations in Appendix C, where we provide in particular the methodology to deduce the effective density of states (eDOS) of surface and bulk magnon-polaritons. The maxima of eDOS indicate that a coherent mode occurs at the given wave vector and frequency. Here, inside the magnet, we consider the physics of the magnon-coupled electromagnetic waves at the surface, with an evanescent wave vector q perpendicular to the slab surface providing a factor of amplitude e^{qz} . The evanescent wave vector should obey the magnetostatic condition that $\nabla \cdot \mathbf{b} = 0$. Meanwhile, we have the continuous condition of magnetic

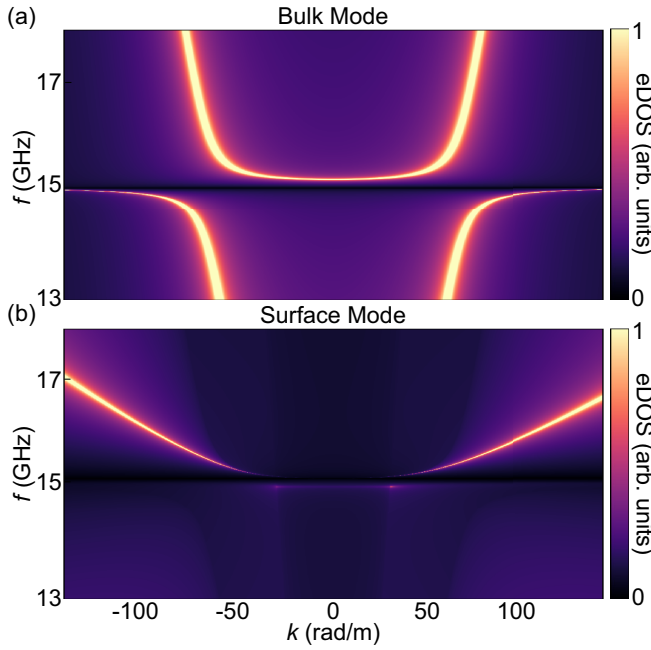


FIG. 4. Reciprocal bulk magnon-polariton and nonreciprocal surface magnon-polariton. (a) The effective DOS (eDOS) of magnon-polariton with a bulk ansatz (see Appendix C), which reproduce the magnon-photon band and its coupling strength in a color scale. The yellow color indicates a mode locating at the given wave vector and the frequency with a similar data process procedure in Ref. [50]. (b) The eDOS of surface magnon-polariton with a finite evanescent wave vector q and an x -direction propagation at the upper surface in a color scale with the same meaning of (a). The surface magnon-polariton appears within the bulk band gap and possesses nonreciprocity, which can be tuned by the external magnetic field.

field at the surface $\partial_{\parallel} \psi_{\text{in}} = \partial_{\parallel} \psi_{\text{out}}$, where $\psi_{\text{in(out)}}$ is the magnetic scalar potential inside (outside) the sample.

As further illustrated in Appendix C, the surface mode dispersion can be deduced using the electric-field boundary condition $e_{y,\text{in}} = e_{y,\text{out}}$, which reads

$$\beta + \frac{q\mu_{33} - ik\mu_{23}}{\mu_{22}\mu_{33} - \mu_{23}\mu_{32}} = 0, \quad (5)$$

where $\beta = \sqrt{k^2 - \omega^2/c^2}$, with c the light velocity in our dielectric media, and μ_{ij} denotes the matrix components of dynamic magnetic permeability $\mu(\omega)$, which can be derived from Eq. (A1), again, whose details are exhibited in Appendix C. Here, β is real since the frequency of surface magnon-polaritons is predicted to appear in the gap, and thus to be lower than the light frequencies. In our calculations, we take the same parameter of other sections, and we choose the propagation direction to be x with an 1.9 GHz coupling gap. If we take $\text{Re}(q) = 0$, we can also deduce the bulk dispersion of magnon-photon bands in hematite, whose details are also exhibited in Appendix C. As illustrated in Fig. 4(a), the well-reproduced bulk dispersion appears at the same frequency of Figs. 1(b) and 1(c), which verifies the legitimacy of our methodology, and further confirms the existence of a band gap of magnon-polaritons in hematite.

Following Eq. (5), with a finite evanescent wave vector q , we obtain the dispersion relation of surface magnon-polaritons in hematite, which are shown by the maxima of eDOS in Fig. 4(b). In agreement with the expectation for the topological bulk band as discussed above, the surface mode appears in the band gap between the topologically trivial and nontrivial bands. In electronic systems, the topological surface (edge) modes tend to connect the valence and conduction bands of electrons. In the magnetic case considered here, the physical interpretation of this surface mode is different. The upper band of the magnon-polariton gradually approaches the light-velocity when k approaches 100 rad/m, which is the largest group velocity of any (quasi)particle, making it impossible for the surface magnon-polariton to connect the “valence” (lower) bulk band and the “conduction” (upper) bulk band. However, our theory indicates that whenever a strongly coupled magnon-photon gap exists, DMI-canted insulating antiferromagnets can host surface magnon-polaritons, enabled by their nontrivial topological nature. Consequently, the surface modes are protected against small perturbations that do not close the gap. Additionally, the group velocity of the surface magnon-polaritons in hematite ($\sim 10^4$ km/s) is still much higher than that of pure magnons (~ 20 km/s) [32–34]. In the polariton regime, the magnetic mode appears as if it “borrowed” some property of the other mode (photons) with which it is in resonance, which has been found in hematite [43]. Likewise, the surface polariton exhibits a long propagation distance, making detection of such modes possible.

One can also observe a nonreciprocal dispersion of the surface mode in Fig. 4(b), which is further illustrated in Appendix D. The nonreciprocal behavior is similar to the nonreciprocity found in Ref. [50], which originates from the magnetic-field-canted sublattice magnetizations with an easy-axis anisotropy. The surface magnon-polariton investigated here is already canted by the DMI and therefore exhibits intrinsic nonreciprocity. We also find that nonreciprocity depends on the sign of the evanescent wave vector q [Eq. (C5)]. This indicates a different propagation preference on either the upper or lower surface. The nonreciprocity of the surface magnon-polariton mode is common in magnonic systems, where time-reversal symmetry is broken and the right-handed precession of dipoles generates a chiral magnetic field [51,54,60]. Unlike strict unidirectional transport of quantized edge modes in electronic systems [71,72] and chiral magnonic bands in magnetic multilayers [53], the surface magnon-polariton mode identified in this study only possesses a nonreciprocal dispersion of different propagation direction $\omega(\mathbf{k}) \neq \omega(-\mathbf{k})$. Such nonreciprocity can be enhanced by the external field, as shown in Appendix D.

V. DISCUSSION AND CONCLUSIONS

In this section, we first give experimental protocols on the excitation and detection of topological surface magnon-polaritons. The excitation of magnons in materials requires wave-vector matching, which is determined by the shape factor of the microwave antenna [8–13]. Here, we propose in addition that the excitation of the topological surface magnon-polariton requires matching of both the in-plane (x - y plane)

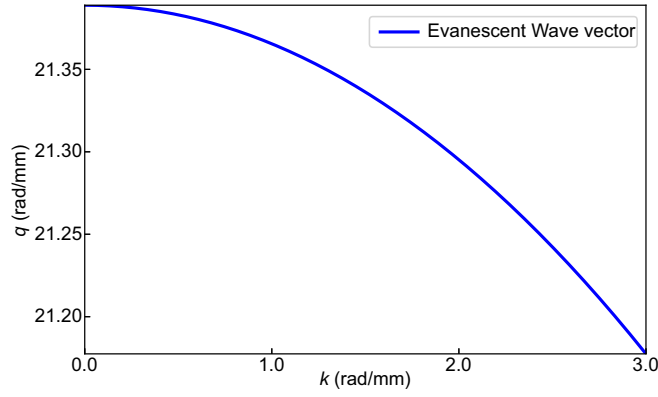


FIG. 5. Matching the evanescent wave vector q . The evanescent wave vector q as a function of the propagating wave vector with $k \parallel \hat{x}$ (both in the unit of rad/mm) for possible experimental realizations of surface magnon-polaritons. The design of microwave antenna should fulfill both the propagating and evanescent wave vectors.

wave vector k and the evanescent wave vector q (along $\hat{z}$). From Eq. (C5), we deduce the evanescent wave vector q in hematite as a function of the in-plane wave vector k , as shown in Fig. 5, where q decreases monotonically with increasing k . In Fig. 5, we use the same parameters as in Fig. 1 and set the propagation direction to be along the x axis. The design of the microwave antenna should explicitly follow the relation shown in Fig. 5. To match k and q , and thus excite magnon-polariton efficiently, a prism with internal reflection can be used [50], where both an in-plane propagating wave vector and an out-of-plane evanescent wave vector can be simultaneously engineered by tuning the antenna shape factor and the total internal reflection angle. We further propose the detection of such modes, which requires time-resolved microwave techniques. A microsecond level resolution is required for a millimeter propagation length to distinguish such mode from bulk modes, after eliminating the signals from cross talks. Applications can be investigated by using the mode to give extra long-range coherent magnon propagation and/or interfere with other modes to provide different magnon precession [36,73–75], where one can refer to the polarization of the surface magnon-polaritons shown in Appendix B. Although the results are deduced in the uses of the parameters from hematite, the physics captured in this work is not sensitive to the exact values of parameters but general to any system with the same low-loss and canting configuration of magnetism. Further investigations can be carried out for a mode with a quantum origin for either a magnon or a photon beyond semiclassical LLM methodology, which can be investigated using the second-quantized approach of a magnon-photon Hamiltonian and quantum transport theories [76,77].

To sum up, we propose an emergent magnon-photon mode, namely topological surface magnon polariton, in an insulating canted antiferromagnet using the parameter of hematite. Our semiclassical model reproduces the strong coupling between magnon and photon modes that occurs in the bulk of such systems by setting a reasonably small net magnetization. The anticrossing band gap between magnon and photon is further confirmed as the origin of the nontrivial Berry curvature of the system. A nonzero Chern number $C = 1$ is found for

the upper bulk band and hosts the existence of a topological surface mode in the system. Adopting our semiclassical model with an evanescent ansatz across the surface, we find the predicted surface mode appears within the bulk band gap with a nonreciprocal feature. Our work extends the methodology to investigate and utilize magnons in antiferromagnetic magnonic devices, and it provides a new functionality for the strong magnon-photon coupling in antiferromagnets.

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DATA AVAILABILITY

The data that support the findings of this article are not publicly available. The data are available from the authors upon reasonable request.

APPENDIX A: ANALYTICAL DETAILS FOR LLM EQUATIONS

In this appendix, we provide details of the Landau-Lifshitz model and the mathematical operations of the Landau-Lifshitz-Maxwell equations for bulk magnon-photon modes. We take into account the free-energy contributions from antiferromagnetic exchange interactions, magnetic anisotropy, DMI, and the Zeeman energy due to a static external magnetic field, following the construction of Ref. [30] (see also [31,32]). In addition, we take into account the Zeeman coupling between the sublattices and the microwave magnetic field, $F_{\text{rf}} = -\mathcal{M}_s(\mathbf{m}_1 + \mathbf{m}_2) \cdot \mathbf{b}$, where $\mathcal{M}_s$ denotes the “saturated” magnetization of each sublattice. Using the total free energy of the two sublattices, the effective fields are obtained as $\mathbf{H}_{\text{eff}}^{1,2} = \frac{\partial F}{\partial (\mathcal{M}_s \mathbf{m}_{1,2})}$. We then further deduce the Landau-Lifshitz equation $\partial \mathbf{m}_{1,2} / \partial t = -\gamma \mu_0 \mathbf{m}_{1,2} \times \mathbf{H}_{\text{eff}}^{1,2}$, as shown in Eq. (1).

To construct LLM equations that are suitable for estimating magnon-photon band structures and characterizing their topological properties, we introduce a transformation from the sublattice magnetizations $\mathbf{m}_{1,2}$ to the net magnetic moment $\mathbf{m} = (\mathbf{m}_1 + \mathbf{m}_2)/2$ and the Néel vector $\mathbf{n} = (\mathbf{m}_1 - \mathbf{m}_2)/2$. As such, the Zeeman coupling term, whose free-energy contribution is proportional to $\mathbf{m} \cdot \mathbf{b}$, is naturally incorporated. Since $\mathbf{m}$ is strictly perpendicular to $\mathbf{n}$, the wave functions defined in this basis are inherently orthogonal, which is rigorous for the calculation of quantum geometric properties. After the basis

transformation, Eq. (1) can be rewritten as

$$\frac{d\mathbf{m}}{dt} = -\gamma\mu_0 \left[\mathbf{m} \times \mathbf{H}_0 + \mathbf{m} \times \frac{\mathbf{b}}{\mu_0} + H_{\text{DM}}((\mathbf{n} \cdot \hat{\mathbf{z}})\mathbf{m} - (\mathbf{m} \cdot \hat{\mathbf{z}})\mathbf{n}) - H_A((\mathbf{m} \cdot \hat{\mathbf{z}})(\mathbf{m} \times \hat{\mathbf{z}}) + (\mathbf{n} \cdot \hat{\mathbf{z}})(\mathbf{n} \times \hat{\mathbf{z}})) + H_a((\mathbf{m} \cdot \hat{\mathbf{y}})(\mathbf{m} \times \hat{\mathbf{y}}) + (\mathbf{n} \cdot \hat{\mathbf{y}})(\mathbf{n} \times \hat{\mathbf{y}})) \right], \quad (\text{A1a})$$

$$\frac{d\mathbf{n}}{dt} = -\gamma\mu_0 \left[\mathbf{n} \times \mathbf{H}_0 + \mathbf{n} \times \frac{\mathbf{b}}{\mu_0} + H_{\text{DM}}((\mathbf{m} \cdot \hat{\mathbf{z}})\mathbf{m} - (\mathbf{n} \cdot \hat{\mathbf{z}})\mathbf{n}) - H_A((\mathbf{m} \cdot \hat{\mathbf{z}})(\mathbf{n} \times \hat{\mathbf{z}}) + (\mathbf{n} \cdot \hat{\mathbf{z}})(\mathbf{m} \times \hat{\mathbf{z}})) + H_a((\mathbf{m} \cdot \hat{\mathbf{y}})(\mathbf{n} \times \hat{\mathbf{y}}) + (\mathbf{n} \cdot \hat{\mathbf{y}})(\mathbf{m} \times \hat{\mathbf{y}})) + 2H_{\text{ex}}(\mathbf{m} \times \mathbf{n}) \right]. \quad (\text{A1b})$$

We then rewrite the Maxwell equations as

$$\frac{\partial \mathbf{b}}{\partial t} = -iK\mathbf{e}, \quad (\text{A2a})$$

$$\frac{\partial \mathbf{e}}{\partial t} = i\frac{K}{\epsilon} \left(\frac{\mathbf{b}}{\mu_0} - M_s \mathbf{m} \right), \quad (\text{A2b})$$

where the matrix K reads

$$K = \begin{pmatrix} 0 & -k_z & k_y \\ k_z & 0 & -k_x \\ -k_y & k_x & 0 \end{pmatrix}. \quad (\text{A3})$$

Alongside Eq. (A1), we then deduce the transformed LLM as the equations of motion of magnon-photon mode in hematite. We rewrite the LLM in a matrix form that reads

$$i\frac{\partial}{\partial t} \mathbf{x}_k = \mathcal{H}_{\text{eff}} \mathbf{x}_k, \quad (\text{A4})$$

where we use the notation of the effective Hamiltonian:

$$\mathcal{H}_{\text{eff}} = \begin{pmatrix} \mathcal{H}_{3 \times 3}^{\text{mag}} & -\theta\gamma\sigma'_2 & O_{3 \times 3} \\ O_{3 \times 3} & O_{3 \times 3} & K_{3 \times 3} \\ \theta\frac{K'}{\epsilon_0} & -\frac{K_{3 \times 3}}{\epsilon_0\mu_0} & O_{3 \times 3} \end{pmatrix}. \quad (\text{A5})$$

In Eq. (A5), we use O to denote the zero matrices, and we use the following definitions:

$$\mathcal{H}_{3 \times 3}^{\text{mag}} = -\gamma\mu_0 \begin{pmatrix} 0 & iH_1 & 0 \\ iH_2 & 0 & iH_3 \\ 0 & iH_4 & 0 \end{pmatrix}, \quad (\text{A6a})$$

$$K' = \begin{pmatrix} -k_z & k_y & 0 \\ 0 & -k_x & 0 \\ k_x & 0 & 0 \end{pmatrix}, \quad (\text{A6b})$$

$$\sigma'_2 = \begin{pmatrix} 0 & 0 & -i \\ 0 & i & 0 \\ 0 & 0 & 0 \end{pmatrix}, \quad (\text{A6c})$$

$$\theta = \frac{H_0 + H_{\text{DM}}}{2H_{\text{ex}}}. \quad (\text{A6d})$$

In Eqs. (A5) and (A6), θ is the canting angle between the Néel vector and the equilibrium value of $\mathbf{m}_1$, which can be found by minimizing the energy of static model, see also [32]. We have

$\theta = \arcsin(H_{\text{DM}} + H_0)/(2H_{\text{ex}}) \ll 1$ since the H_0 and H_{DM} is much smaller than H_{ex} in hematite. However, if one discusses such a model in other materials, this approximation should be reexamined. For the definitions in Eq. (A6), we have

$$H_1 = \frac{H_0 + H_{\text{DM}}}{2H_{\text{ex}}} H_A + H_0 - H_{\text{DM}}, \quad (\text{A7a})$$

$$H_2 = \frac{H_0(H_A - 2H_{\text{ex}}) + H_A H_{\text{DM}}}{2H_{\text{ex}}}, \quad (\text{A7b})$$

$$H_3 = H_a, \quad (\text{A7c})$$

$$H_4 = \frac{(H_0 + H_{\text{DM}})H_{\text{DM}} - 2H_{\text{ex}}(H_a + H_A + 2H_{\text{ex}})}{2H_{\text{ex}}}. \quad (\text{A7d})$$

APPENDIX B: WAVE-VECTOR-DEPENDENT POLARIZATION OF BULK AND SURFACE MAGNON-POLARITONS

In this appendix, we show the characters of the net magnetic moment precession polarization. Figure 6(a) shows two curves: the upper curve represents the lower magnon, which decreases monotonically from 6.5 to 5, while the lower curve, corresponding to the upper band, shows a monotonic decrease starting from 5. Similar to the dispersion curves, an anticrossing phenomenon is observable in this figure. The physical origin of this anticrossing lies in the interaction between magnons and photons. In hematite, the anisotropic magnon-photon coupling, described by LLM equations, plays a crucial

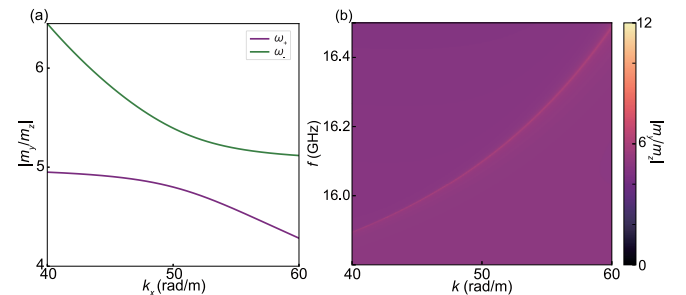


FIG. 6. Wave-vector-dependent polarization of bulk and surface magnon-polaritons. (a) The polarization $|m_y/m_z|$ of upper (ω_+) and lower (ω_-) branches of bulk magnon-photon band illustrated by the purple and green curve, respectively. (b) The polarization of surface magnon-polaritons sketched by the colorbar.

role. At the anticrossing point, strong magnon-photon hybridization occurs, leading to a non-negligible out-of-plane component ($\mathbf{b}_z$) of the photon's oscillating magnetic field, which affects the magnetic moment precession, altering the ratio m_y/m_z . Specifically, in the anticrossing region, the m_y/m_z is approximately 6, which corresponds to the fact that the band gap in the k_y direction is six times that in the $k_x(k_z)$ direction. Increasing k_x shifts the mode toward magnon dominance, the DMI stabilizes this regime, and, as a result, the ratio decreases. Rather than suppressing m_z , it shifts the relative amplitudes of m_y and m_z , so the m_y/m_z ratio decreases and deviates from 6. This behavior validates the conclusion that anisotropic coupling strength stems from the precession geometry of $\mathbf{m}$.

Figure 6(b) is a heatmap of m_y/m_z in the frequency-wave-vector dispersion plane, covering the energy range of the bulk magnon-photon gap and adjacent bulk modes. Two distinct regions are clearly distinguishable in the figure: the first is the bulk mode region, encompassing areas with lower frequencies (lower bulk band) and higher frequencies (upper bulk band), where m_y/m_z exhibits uniformly low values (2–5) with minimal variation across k_x . The second is the bulk-band-gap region, which contains a narrow bright yellow strip. The trajectory of this strip in the plane perfectly matches the dispersion relation of the topological surface magnon-polariton, derived from Eq. (5) in the main text, and within the strip, it stabilizes at approximately 10. The significant difference in the ratio between surface and bulk modes arises from the topological origin and interface-bound nature of the surface magnon-polariton: the vacuum environment lacks a magnetic medium to support m_z -coupled oscillations, leading to the suppression of m_z . In contrast, m_y can couple to the evanescent wave of photons at the interface—sustained by the dielectric response of hematite, resulting in the enhancement of m_y and the final ratio of $m_y/m_z \sim 10$.

APPENDIX C: DETAILS FOR SURFACE-MODE CALCULATIONS

In this appendix, we exhibit the details for the calculations of surface magnon-polariton mode. Our method follows explicitly Ref. [50]. First, we give the magnetic susceptibility $\chi(\omega)$. In Eq. (A1), we have $\mathbf{m}_{1,2}(\omega) = \frac{1}{i\omega} \mathbf{m}_{1,2}(\omega) \times \mathbf{H}_{\text{eff}}^{1,2}$, which can be rewritten as

$$\begin{pmatrix} m_{1,2}^y \\ m_{1,2}^z \end{pmatrix} = \begin{pmatrix} \chi_{1,2}^{yy} & \chi_{1,2}^{yz} \\ \chi_{1,2}^{zy} & \chi_{1,2}^{zz} \end{pmatrix} \begin{pmatrix} h^y \\ h^z \end{pmatrix}, \quad (\text{C1})$$

where we use the intrinsic frame of references for $\mathbf{m}_{1,2}$. We then transform the susceptibility to the $\mathbf{m}, \mathbf{n}$ basis, which reads

$$\begin{aligned} m^y &= (m_1^y - m_2^y) \cos \theta, \\ m^z &= m_1^z + m_2^z, \end{aligned} \quad (\text{C2})$$

where we can derive the total magnetic susceptibility $\mathbf{m} = \chi \mathbf{h}$. Thus, we have the total permeability,

$$\mu(\omega) = \begin{pmatrix} \mu_{11} & 0 & 0 \\ 0 & \mu_{22} & \mu_{23} \\ 0 & \mu_{32} & \mu_{33} \end{pmatrix} = \text{diag}(\mu_{11}, I_2 + \mu_0 \chi), \quad (\text{C3})$$

which is useful in the derivation of the dispersion. Here we note that μ_{11} is irrelevant to the calculation of dynamic modes in our work, and I_2 denotes the unit matrix. For the surface mode of magnon-polaritons in insulators, we have the ansatz of the magnetic scalar vector,

$$\psi_{\text{in}} = \psi_{0,\text{in}} e^{qz} e^{i(kx - \omega t)}, \quad (\text{C4a})$$

$$\psi_{\text{out}} = \psi_{0,\text{out}} e^{-\beta z} e^{i(kx - \omega t)}, \quad (\text{C4b})$$

where we assume the propagation direction is x , and we are considering the upper surface of the slab. $\nabla \cdot \mathbf{B} = 0$ in the vacuum gives $\beta = \sqrt{k^2 - \frac{\omega^2}{c^2}}$ in the ansatz. At the surface, we have two continuity conditions, namely that the electric displacement vector normal to the surface and the magnetic field parallel to the surface are continuous. In addition, in Maxwell's equations, we have $\nabla \times \mathbf{e} = -\partial \mathbf{b} / \partial t$ at the surface, which gives $e_{y,\text{in}} = e_{y,\text{out}}$. Substituting the ansatz to the boundary conditions, we obtain [50]

$$\frac{\omega^2}{c^2 \beta} \frac{-q\mu_{33}(\omega) + ik\mu_{23}(\omega)}{\mu_{22}(\omega)\mu_{33}(\omega) - \mu_{23}(\omega)\mu_{32}(\omega)} = 1, \quad (\text{C5})$$

which can be further simplified as Eq. (5), and the left-hand side is the effective density of states of the surface magnon-polaritons. eDOS goes to 1 represents the appearance of the coherent mode [50]. If we substitute the plane-wave ansatz $\psi_{\text{in}} = \psi_{0,\text{in}} e^{qz} e^{i(kx - \omega t)}$ with Maxwell's equation inside the bulk, we obtain

$$\frac{1}{\epsilon\omega^2} \frac{k^2 \mu_{33}(\omega)}{\mu_{22}(\omega)\mu_{33}(\omega) - \mu_{23}(\omega)\mu_{32}(\omega)} = 1, \quad (\text{C6})$$

whose left-hand side is also a function of ω and k . If we plot the left-hand side, the maxima which hit the quantity 1 give the solution to the bulk magnon-photon mode.

APPENDIX D: MAGNETIC FIELD TUNABILITY OF NONRECIPROCAL SURFACE MAGNON-POLARITONS

This appendix supplements the modulation of topological surface states by external magnetic fields, focusing on the magnetic field response of nonreciprocal surface magnon-polaritons. Figure 7 shows the effective density of states (eDOS) of surface magnon-polaritons under three distinct external magnetic field strengths ($H_0 = 60$ mT, $H_0 = 500$ mT, $H_0 = 1$ T), consistent with the experimental field range specified in the main text. In these eDOS maps, peak intensity characterizes the coherent modes, meanwhile, the asymmetry related to the direction of in-plane wave vector quantifies the degree of nonreciprocity (a core feature of the topological surface modes discussed in the main text). The two-sublattice system is fundamentally characterized by the DMI, which tilts the magnetization of both sublattices. This tilted state imparts intrinsic nonreciprocal behavior to the system. Notably, this nonreciprocal characteristic exhibits high sensitivity to variations in the external magnetic field H_0 . Direct observations from Fig. 7 (corresponding to the three aforementioned values) reveal a clear trend: as H_0 increases, the asymmetry of the eDOS peaks becomes increasingly pronounced. Specifically, the eDOS peaks at $H_0 = 60$ mT show relatively weak directional asymmetry; at $H_0 = 500$ mT, the peaks exhibit a more significant bias toward one propagation

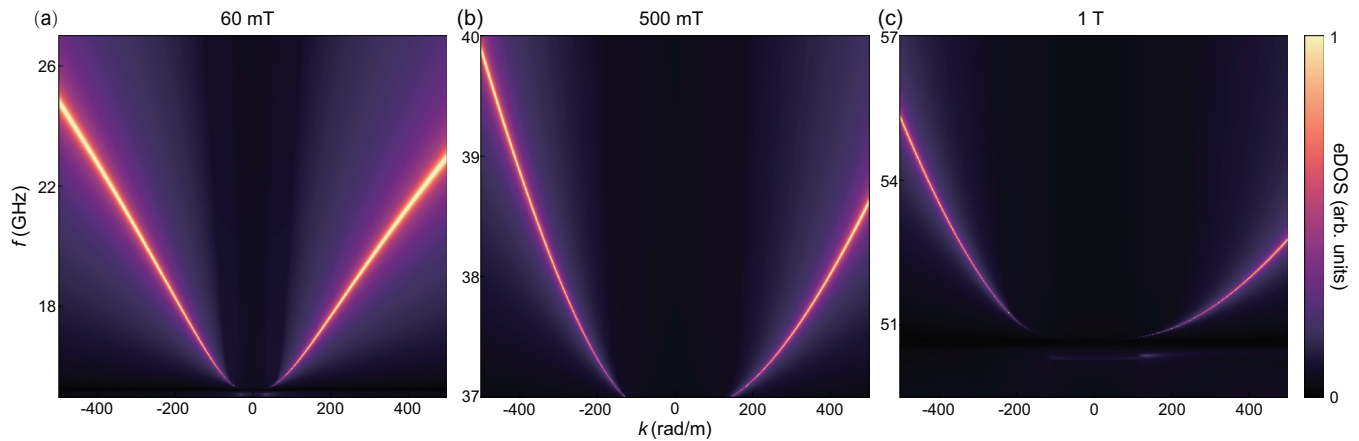


FIG. 7. Magnetic field tunability of nonreciprocal surface magnon-polaritons. Effective density of states (eDOS) of surface magnon-polaritons under different external magnetic field configurations: (a) for 50 mT, (b) for 500 mT, and (c) for 1 T, demonstrating the nonreciprocity of surface magnon-polaritons is able to be tuned by the external magnetic field.

direction; and at $H_0 = 1$ T, this asymmetric feature is further enhanced, with the eDOS peak asymmetry reaching its maximum within the measured field range.

The physical origin of this trend lies in the modulation of sublattice magnetization tilted by DMI and external field. As H_0 increases, it combines with the intrinsic H_{DM} , further amplifying the asymmetry of the magnetization distribution. This more pronounced magnetization asymmetry, in turn, enhances the directional difference in the coupling between

surface magnon-polaritons and the underlying magnetic order: the mode propagates more strongly in one direction (corresponding to higher frequency) and more weakly in the opposite direction (corresponding to lower frequency). This mechanism directly explains the increasingly asymmetric eDOS peaks in Fig. 7 and provides experimental support for the main text's conclusion that nonreciprocal topological surface modes can be tuned by external magnetic fields.

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