Magnon waveguide with nanoscale confinement constructed from topological magnon insulators

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Topological magnon insulators host spatially confined edge magnons brought about by the Dzyaloshinskii-Moriya interaction. Bringing two topological magnon insulators into contact results in topologically protected unidirectional interface magnons. These interface modes decay rapidly toward the bulk regions of the sample. As a result, heat and spin currents associated with these magnons are as well unidirectional and strongly confined to a few-nanometer-wide strip along the interface. On top of this, these interface currents follow any geometry owing to the topological nature of the magnons. In this theoretical study, we propose and analyze two recipes for composing magnon waveguides with nanoscale confinement, one from topologically different phases, another from identical phases. We further identify material classes to construct these magnon waveguides and propose an experiment to verify their topological nature.

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Introduction. Lu2V2O7, an insulating ferromagnetic pyrochlore oxide, exhibits the magnon Hall effect [1], i.e., a transverse heat current upon application of a longitudinal temperature gradient. This transverse heat current is due to uncompensated edge currents carried by magnons [2,3]. The edge currents originate from the nontrivial topology of the magnon band structure of Lu2V2O7 and are mathematically described in terms of a Berry curvature [4]. The nontriviality is a result of the Dzyaloshinskii-Moriya (DM) interaction that is brought about by spin-orbit interaction and the lack of inversion symmetry of the pyrochlore lattice. In analogy to electronic topological insulators, Lu2V2O7 is a “topological inversion symmetry of the pyrochlore lattice. In analogy to...
propagation direction of the edge states depends solely on the edge orientation.

Interface magnon modes are investigated for two semi-infinite samples sharing a common interface (as shown in Fig. 1). The magnon band structure is analyzed in terms of the spectral density computed by Green-function renormalization [10,18,19].

The heat current associated with each magnon is computed for a slab geometry, that is, vacuum/phase (dashed line) along the y direction. As a result, a topological interface magnon moves unidirectionally along this interface (blue arrow). It decays rapidly toward the two bulk regions of the sample, i.e., in both positive and negative x direction. Consequently, the heat current \( j_y^x(x) \) associated with this interface magnon is strongly confined to a few-nanometer-wide strip along the interface.

Results and discussion. We demonstrate our prediction by means of a kagome slab constructed from two topologically different phases A and B with Chern triples \((-1,0,1)\) and \((-1,2,-1)\), respectively, each occupying 48 atomic layers (Fig. 2). We performed a considerable number of calculations for various sets of exchange parameters, topological phases, and slab widths \( L \), all exhibiting the general features discussed in the following.

The heat current profile of the slab shows three eye-catching features [Fig. 2(a)]. At the left (right) edge the current is negative (positive), is maximum (in absolute value) at the outermost layers, and drops to (almost) zero within about ten atomic layers. This spatial dependence is readily explained by the decay of the magnon edge states toward the bulk [10], in accordance with Ref. [5]. Note that these edge currents are an artifact of the slab geometry; they do not appear in a true semi-infinite system.

We recall that (i) a positive (negative) winding number indicates current in the positive (negative) y direction and (ii) topological edge magnons located at opposite edges of a single-phase slab propagate in opposite directions. Within this respect it is important to note that the band structures of A [Fig. 2(b)] and B [Fig. 2(d)] are shown for opposite edges.

As a result for the two-phase slab, the left edge current has negative sign, because the winding numbers of A read \( \nu_1^A = \nu_2^A = -1 \). In contrast, there are both left and right moving magnon states at the edge of phase B as its winding numbers differ in sign, \( \nu_1^B = -1 \) and \( \nu_2^B = +1 \). The edge mode lower in energy and with positive slope is more occupied than the descending one at a higher energy; consequently, it dominates the edge current which is positive. The black arrows in (a) represent the number and propagation direction of the edge modes which follows from the spectral densities shown in the insets (b) and (d). The edge modes are easily identified: they cross the bulk band gaps and connect adjacent bulk bands.

The exchange parameters \( J_N \), \( J_{NN} \), and \( D \) of phase A—chosen to mimic Lu\(_2\)V\(_2\)O\(_7\)—allow us to estimate the order of magnitude of the local heat current density. An analysis of inelastic neutron-scattering experiments probing the spin dynamics in Lu\(_2\)V\(_2\)O\(_7\) gave a \( D/J_N \) ratio of 0.18 with \( J_N = 8.22(2) \) meV [22]. Furthermore, it confirmed that nearest-neighbor exchange dominates. However, from the transversal thermal conductivity \( \kappa_{xy} \) of the magnon Hall effect one obtains a \( D/J_N \) ratio from 0.32 [1] to 0.38 [23]. Additionally, density functional theory yielded \( D/J_N = 0.05 \) with \( J_N = 7.09 \) meV for Y\(_2\)V\(_2\)O\(_7\) that possesses magnetic properties similar to Lu\(_2\)V\(_2\)O\(_7\) [24]. Considering the above numbers we chose \( J_N = 7.5 \) meV, \( D = 1.8 \) meV, and \( J_{NN} = 0 \) meV for phase A. Again, we point out that qualitative results do not depend on this particular choice.

The exchange parameters of phase B were deliberately chosen such that it is in a different topological phase and that band gaps of A and B coincide [Figs. 2(b)–2(d)]. Thus, the interface winding numbers \( \nu_i^B = \nu_i^A \) for \( i = 1, 2, 3 \) and \( n,m = 1,2 \) are uniquely defined within mutual band gaps of phases A and B. The winding numbers of B have to be subtracted from those of A because their interface normal vectors point in opposite directions (A: right edge; B: left edge). This definition is fully in line with that given in

FIG. 1. (Color online) Sketch of a magnon waveguide. Two kagome ferromagnets, phases A (green) and B (red), share a common interface (dashed line) along the y direction. As a result, a topological interface magnon moves unidirectionally along this interface (blue arrow). It decays rapidly toward the two bulk regions of the sample, i.e., in both positive and negative x direction. Consequently, the heat current \( j_y^x(x) \) associated with this interface magnon is strongly confined to a few-nanometer-wide strip along the interface.

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The interface winding numbers tell the number and propagation direction of topological interface magnons, which is confirmed numerically for a manifold of topologically different kagome interfaces.

Due to the energy positions of the band gaps of A and B, there are three interface band gaps [Fig. 2(c)]. There is no topological magnon mode inside the lowest interface band gap since $v_0^A - v_0^B = (-1) - (-1) = 0$. For the second interface band gap one has to subtract the second winding number of B from the first winding number of A, $v_3^A - v_3^B = (-1) + (+1) = 2$; hence, there are two topological magnon modes. Eventually, the second band gap of A has no partner in B, yielding $v_4^A = v_4^B = -1$. As a consequence, there are in total three magnon interface modes $(-2 - 1) = 3$ propagating in the positive y direction [black arrows in Fig. 2(a) and spectral density in Fig. 2(c)]. They transport heat unidirectionally with a local heat current density of about $10^{-3}$ W/m. The current profile is maximum at the interface, i.e., about layer 48, and is similar to that in the sketch of the magnon waveguide (Fig. 1).

To demonstrate that the origin of the interface current is the topological interface modes we analyze the energy-resolved contribution $j_y(e,x)$ to the local current density, both around the interface, i.e., between atomic layers 45–50, and within phase B, i.e., between atomic layers 65–70; cf. the red regions in the upper panels of Figs. 3(a) and 3(b). At the interface the energy-resolved current shows large contributions within the energy range of the interface modes and is small elsewhere [lower panel of Fig. 3(a)]. The bulk bands provide magnons propagating in both directions, which results in a zero net current. In contrast, within phase B the energy-resolved current is small over the entire energy range of the magnon dispersion because there are no topological interface modes in these layers [lower panel of Fig. 3(b)].

The chiral interface modes are the key ingredient for the magnon waveguide. The heat current profile in Fig. 2(a) shows a large peak at the interface. It is confined to a strip approximately 20 atomic layers wide, i.e., $L_{\text{conf}} \approx 5a$ (a is the lattice constant) for the geometry shown in Fig. 1. For a kagome lattice constant $a = 7.024$ Å associated with the vanadium sublattice of Lu$_2$V$_2$O$_7$ [1], $L_{\text{conf}}$ is smaller than 4 nm. Being topologically protected modes, they unconditionally follow any geometry of the interface, e.g., corners.

For an experimental realization of the above magnon waveguide one needs to identify samples in different topological phases. One family of crystals to consider are insulating (or semiconducting) ferromagnetic pyrochlore oxides as these show the magnon Hall effect [1,23]. The sign of the transversal thermal conductivity $\kappa_{xy}$ is intimately related to the topological phase of the crystal which also dictates the propagation directions of the edge modes [10]. Exhibiting a negative $\kappa_{xy}$ in magnon Hall experiments, In$_2$Mn$_2$O$_7$ is very likely in a different topological phase than Lu$_2$V$_2$O$_7$ and Ho$_2$V$_2$O$_7$, both showing positive $\kappa_{xy}$ [23]. From a fit of the thermal Hall conductivity the authors of Ref. [23] concluded that the DM constant of In$_2$Mn$_2$O$_7$ differs in sign from those of the other two representatives. A sign change of $D$ causes a sign change of the Chern numbers, which puts In$_2$Mn$_2$O$_7$ into a different topological phase. Thus, an interface built from In$_2$Mn$_2$O$_7$ and, say, Lu$_2$V$_2$O$_7$—with a lattice mismatch of about 2.3% (Refs. [1,26])—is very likely to host topological interface magnons and, thus, would show magnon waveguide behavior.

Having established magnon waveguides at interfaces of topologically different materials, we now discuss briefly that even at interfaces of topologically identical materials these modes exist. Being bosons, all magnons contribute to the heat current, in contrast to electrons which contribute at the Fermi energy only. Hence, the bulk-boundary correspondence has to be reviewed for all band gaps.
neighbor interaction and heat current at their interface. Upon increasing the nearest-neighbor interaction, no topologically nontrivial interface states and, thus, no local gaps match and their winding numbers are equal, there are no topological phase transitions. Since \( \nu_B = -1 \) (the winding number of phase \( B \)), it stays in its topological phase \( \nu = 1 \) for the topmost interface magnon. By heating locally, for example, by a laser pulse, both bulk and interface magnons become occupied. By measuring thermographic profiles along the interfaces shortly after the heating (i.e., in a “pump and probe” experiment) trivial and nontrivial samples can be distinguished: the thermographic profile caused by the isotropic heat transport of the bulk magnons is determined by the longitudinal thermal conductivities \( \kappa_{xx} \) of phases \( A \) and \( B \), respectively, that is, it is radially symmetric within each phase. Trivial interface magnons cause a symmetric heat transport along the interface, whereas the (unidirectional) nontrivial magnons result in an asymmetric heat transport that is determined by the transverse thermal conductivity \( \kappa_{xy} \). Thus, the topologically nontrivial nature of these interface magnons manifests itself by a deformed thermographic profile. The magnitude of this deformation is given by the Hall angle \( \kappa_{xy}/\kappa_{xx} \), which is of the order of \( 10^{-3} \) for \( \text{Lu}_2\text{V}_2\text{O}_7 \) [1].

Terahertz radiation [28] could be utilized for transport in the proposed waveguide. At low temperatures mainly the lowest magnon band is occupied. Thus, incident radiation with a suitable energy range will excite preferably the topological interface magnons, with the consequence that the associated spin and heat current propagates unidirectionally along the interface.

FIG. 3. (Color online) Energy-resolved contribution to the local heat current density. Upper panels show the current profile for the slab in Fig. 2. The layers marked in red around the interface [(a), atomic layers 45–50] and within the \( B \) region [(b), atomic layers 65–70] were considered for the energy-resolved current. Lower panels show the spectral density and the energy-resolved current density \( J_q^y(\varepsilon,x) \) of said layers. Temperature \( T = 40 \text{ K} \), energy interval 0.52 meV. The “wriggling” of \( J_q^y(\varepsilon,x) \) about 0 is attributed to numerics.

Consider two identical crystals \( A \) and \( B \) both belonging to the topological phase, say \((-1,0,1)\). Because their band gaps match and their winding numbers are equal, there are no topologically nontrivial interface states and, thus, no local heat current at their interface. Upon increasing the nearest-neighbor interaction \( J_N \) of \( B \), it stays in its topological phase \( B \), however, the total width of its band structure is increased. More precisely, the uppermost band of \( B \) entirely tops the uppermost band of \( A \), which introduces a new band gap at the interface. The winding number of this band gap is determined solely by \( v_2^B = -1 \) (the winding number of phase \( A \) equals zero at this energy because the sum of all Chern numbers equals zero). Since \( v_1^B = -v_2^B = +1 \) for the topmost interface band gap, there exists one interface state although both \( A \) and \( B \) belong to the same topological phase [27].

An advantage of the isophase waveguide is that the exchange parameters could be tuned by doping one of the materials. A shortcoming seems to be related to the energy region of the interface mode: an elevated energy leads to both a small occupation and a small heat current at a given temperature.

Concluding remarks. To prove experimentally topological properties of interface magnons we suggest to make use of the unidirectional heat current at the interface. Consider two samples, one that does and another that does not host a topological interface magnon. By heating locally, for example by a laser pulse, both bulk and interface magnons become occupied. By measuring thermographic profiles along the interfaces shortly after the heating (i.e., in a “pump and probe” experiment) trivial and nontrivial samples can be distinguished: the thermographic profile caused by the isotropic heat transport of the bulk magnons is determined by the longitudinal thermal conductivities \( \kappa_{xx} \) of phases \( A \) and \( B \), respectively, that is, it is radially symmetric within each phase. Trivial interface magnons cause a symmetric heat transport along the interface, whereas the (unidirectional) nontrivial magnons result in an asymmetric heat transport that is determined by the transverse thermal conductivity \( \kappa_{xy} \). Thus, the topologically nontrivial nature of these interface magnons manifests itself by a deformed thermographic profile. The magnitude of this deformation is given by the Hall angle \( \kappa_{xy}/\kappa_{xx} \), which is of the order of \( 10^{-3} \) for \( \text{Lu}_2\text{V}_2\text{O}_7 \) [1].

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Concerning an experimental realization, one family of choice is ferromagnetic insulating pyrochlore oxides: the vanadates \( X_2\text{V}_2\text{O}_7 \) (\( X = \text{Lu}, \text{Yb}, \text{Tm} \)) with a Curie temperature of about 70 K and manganates \( X_2\text{Mn}_2\text{O}_7 \) (\( X = \text{Ho}, \text{Er}, \text{Yb}, \text{Lu}, \text{Y} \)) with Curie temperatures in the range 20–40 K [26]. In addition, the insulating ferromagnetic spinels are promising candidates; for example, the chromium spinels \( X\text{Cr}_2\text{Y}_4 \) (\( X = \text{Cd}, \text{Hg} \) and \( Y = \text{S}, \text{Se} \)) show strong first-neighbor interactions \( J_N \) that could fulfill the aforementioned requirements. \textit{Ab initio} calculations yield a range of exchange parameters allowing for a factor of about \( J_N^{1B}/J_N^{1B} \approx 1.16 \) [29]. Apart from pyrochlore crystals, distorted perovskites, like BiMnO\(_3\), come into question, as they exhibit the magnon Hall effect, too [23].

Finally, we note that the organometallic magnet \( \text{Cu}(1,3\text{-benzenedicarboxylate}) \) provides stacked kagome lattices [30] with very small interlayer interaction. Ferromagnetic order is present below 1.77 K with an intraplane interaction ratio of \( J_B/J_{1B} \approx 15\% \) (Ref. [31]). Thus, \( \text{Cu}(1,3\text{-benzenedicarboxylate}) \) is a topological magnon insulator.

Our proposal of magnon waveguides built from topological magnon insulators calls for \textit{ab initio} electronic structure calculations for determination of their exchange and Dzyaloshinskii-Moriya parameters as well as their topological phase.

Unidirectional magnon edge states offer a way to control topological excitations, e.g., skyrmions, thereby showing potential for applications in spintronics as investigated in Ref. [32]. The chiral interface states extend possibilities of manipulation of said local topological excitations.

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[6] Note the difference to “topological magnonic crystals.” Their nontriviality is due to the dipolar interaction, i.e., they provide chiral edge modes for magnetostatic spin waves [7].