

Giant Uncompensated Magnon Spin Currents in X-Type Magnets

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Magnon spin currents in insulating magnets are useful for low-power spintronics. However, in magnets stacked by antiferromagnetic (AFM) exchange coupling, which have recently aroused significant interest for potential applications in spintronics, Bose–Einstein distribution populates magnon states across all energies from opposite eigenmodes, and hence the magnon spin current is largely compensated. Contrary to this common observation, here, we show that magnets with X-type AFM stacking, where opposite magnetic sublattices form orthogonal intersecting chains, support giant magnon spin currents with minimal compensation. Our model Hamiltonian calculations predict magnetic chain locking of magnon spin currents in these X-type magnets, significantly reducing their compensation ratio. In addition, the one-dimensional nature of the chain-like magnetic sublattices enhances magnon spin conductivities surpassing those of two-dimensional ferromagnets and canonical altermagnets. Notably, uncompensated X-type magnets, such as odd-layer antiferromagnets and ferrimagnets, can exhibit magnon spin currents polarized opposite to those expected by their net magnetization. These unprecedented properties of X-type magnets, combined with their inherent advantages resulting from AFM coupling, offer a promising new path for low-power high-performance spintronics.

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Spin currents play the key role in spintronics for information read-out and write-in.^[1–8] Conventionally, spin current generation is associated with the electric-field-driven motion of electrons which is spin polarized by exchange or spin-orbit interactions. In contrast to electrons carrying both charge and spin, magnons—quantized spin waves, only carry spin angular momentum and thus can form a pure spin current without an accompanying electric current.^[9–14] In this regard, magnons have obvious advantages over electrons, especially due to their possible generation in magnetic insulators which substantially reduces energy dissipation.^[10,11]

The magnon spin current can be generated by sev-

eral external stimuli, such as a temperature gradient,^[15,16] a microwave,^[17] a photoexcitation,^[18,19] etc. These approaches have been mostly demonstrated for ferromagnetic (FM) or ferrimagnetic (FIM) materials due to their spontaneous magnetization, while the efficiency of magnon spin current generation is yet to be improved. Recently, giant magnon spin conductivity has been reported for ultrathin FIM films supported by the two-dimensional (2D) confinement.^[20,21] This achievement suggests an efficient route for high-performance magnonic applications using 2D systems such as ultrathin films and/or van der Waals materials. The robustness of the magnetic order in the 2D limit and interface (surface) scattering^[22] are the existing

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challenges of this approach.

Recently, magnets stacked by antiferromagnetic (AFM) exchange coupling have attracted increasing interest for spintronics due to the spin dynamics being much faster and the stray fields being much smaller than those in FM systems.^[23–27] Magnon excitations in AFM-stacked magnets typically occur at much higher frequencies, often in the terahertz range, compared to ferromagnets. However, due to the counteraction of the magnetic sublattices, the magnon spin currents in AFM-stacked magnets are expected to be strong only for ferrimagnets with large net magnetization. For example, a spin Seebeck effect in conventional antiferromagnets can be induced under an applied magnetic field, which only marginally breaks magnetic compensation.^[28–32] A spin Nernst effect may occur even when the chiral modes are degenerate.^[33–36] However, this effect primarily originates from the inter-band contributions of certain specific states, which cannot guarantee sizable magnon spin conductivity. Unconventional antiferromagnets, including altermagnets^[37–48] and some non-collinear antiferromagnets,^[49–56] have been proposed to host non-relativistically spin-split band structures supportive of a highly spin-polarized^[49,57] or even half-metallic^[58,59] Fermi surface capable of generating a strong spin current. However, although these materials support spin-split magnon bands,^[60–66] this property is not sufficient for strong magnon spin currents. This is because the Bose–Einstein distribution of magnons gathers the contribution from states across all energies of the opposite magnetic sublattices, and hence the magnon spin current might still be largely compensated.

In this Letter, we show that the efficient generation of magnon spin currents is possible by using the recently discovered antiferromagnets and ferrimagnets with X-type AFM stacking.^[67] This stacking employs one-dimensional (1D) FM-ordered atomic chains with opposite magnetic moment orientations (denoted as A- and B-chains) as the fundamental building blocks. Instead of alternating chains parallel to each other as in C-type stacking, the X-type stacking hosts alternating layers of A- and B-chains parallel to each other within the plane but crossing in the adjacent layers. Due to this unique cross-chain pattern of the two magnetic sublattices, these X-type AFM-stacked magnets (or X-type magnets for short) support barely compensated magnon spin currents with giant magnon spin conductivities. These results open a new route for efficient magnon spin-current generation for low-power spintronics.

We start from considering a collinear AFM-stacked magnet with two magnetic sublattices that is described by the Heisenberg Hamiltonian,^[68]

$$\begin{aligned}
 H = & -J_{\text{FA}} \sum_i \mathbf{M}_{\text{A}_i} \cdot \mathbf{M}_{\text{A}_{i+1}} - J_{\text{FB}} \sum_i \mathbf{M}_{\text{B}_i} \cdot \mathbf{M}_{\text{B}_{i+1}} \\
 & - J_{\text{AF}} \sum_{i,j} \mathbf{M}_{\text{A}_i} \cdot \mathbf{M}_{\text{B}_j} + K_{\text{A}} \sum_i (\hat{\mathbf{n}} \cdot \mathbf{M}_{\text{A}_i})^2 \\
 & + K_{\text{B}} \sum_i (\hat{\mathbf{n}} \cdot \mathbf{M}_{\text{B}_i})^2,
 \end{aligned} \quad (1)$$

where the local moment $\mathbf{M}_{\text{A}}$ ($\mathbf{M}_{\text{B}}$) in sublattice A (B) is indexed by i or j , J_{FA} (J_{FB}) > 0 is the nearest-neighbor

FM exchange parameter in sublattice A (B), $J_{\text{AF}} < 0$ is the nearest-neighbor inter-sublattice AFM exchange parameter, $\hat{\mathbf{n}}$ is the Néel vector, and K_{A} (K_{B}) is the magnetic anisotropy of sublattice A (B). A temperature gradient $\nabla_m T$ along direction m drives a magnon spin current

$$\mathbf{j}_n^s = \sigma_{nm}^s \nabla_m T. \quad (2)$$

In AFM-stacked collinear magnets, there are two eigenmodes of magnons with opposite chiralities,^[68] conventionally referred to as the α and β modes. Due to their opposite chiralities, these two modes carry different spins during propagation. Therefore, the thermal magnon spin conductivity σ_{nm}^s can be decomposed into two eigenmodes, α and β , so that

$$\sigma_{nm}^s = \sigma_{nm}^\alpha + \sigma_{nm}^\beta. \quad (3)$$

For the longitudinal ($n = m$) magnon spin conductivity generated by the spin Seebeck effect, we have^[60]

$$\begin{aligned}
 \sigma_{nn}^i &= \int \frac{d^2 \mathbf{k}}{(2\pi)^2} \sigma_{nn}^i(\mathbf{k}) \\
 &= \frac{\gamma_0^i \tau_0 \hbar}{k_{\text{B}} T^2} \int \frac{d^2 \mathbf{k}}{(2\pi)^2} [v_n^i(\mathbf{k})]^2 E^i(\mathbf{k}) f[E^i(\mathbf{k})],
 \end{aligned} \quad (4)$$

where $i \in \{\alpha, \beta\}$, $\gamma_0^\alpha = 1$, $\gamma_0^\beta = -1$, τ_0 is the relaxation time of magnon, $E^{\alpha(\beta)}(\mathbf{k})$ is the magnon band energy at wave vector $\mathbf{k}$ contributed by eigenmode α (β), $v_n^{\alpha(\beta)}$ is the group velocity along direction n , T is temperature, and $f(E) = \frac{e^{E/k_{\text{B}}T}}{(e^{E/k_{\text{B}}T} - 1)^2}$.

Since σ_{nn}^α and σ_{nn}^β have opposite signs, a ratio defined by $\xi = \frac{\min\{|\sigma_{nn}^\alpha|, |\sigma_{nn}^\beta|\}}{\max\{|\sigma_{nn}^\alpha|, |\sigma_{nn}^\beta|\}}$ can be regarded as a measure of the magnon spin current compensation by the opposite magnetic sublattices, where $\xi = 0$ implies no compensation, while $\xi = 100\%$ implies full compensation so that $\sigma_{nn}^s = 0$. For example, conventional collinear antiferromagnets with $\mathbf{M}_{\text{A}} = -\mathbf{M}_{\text{B}}$ and degenerate magnon bands exhibit perfect compensation ($\xi = 100\%$) and thus vanishing σ_{nn}^s . In unconventional collinear antiferromagnets, including altermagnets^[37–47] and other non-relativistically spin-split antiferromagnets,^[59] magnon bands are non-degenerate, allowing for $\xi < 100\%$, and thus for a finite σ_{nn}^s .^[60–62] However, due to the Bose–Einstein distribution of magnons collecting contributions from all energies [Eq. (4)] including those with zero spin splitting, the resulting ξ may still be close to 100%, thus hindering a sizable σ_{nn}^s . This is also the case for the uncompensated AFM-stacked magnets with $\mathbf{M}_{\text{A}} \neq -\mathbf{M}_{\text{B}}$; i.e., ferrimagnets.

Now, we predict that the X-type AFM stacking can effectively suppress ξ , making the thermal magnon spin current barely compensated. The X-type AFM stacking is an extension of the A-, C-, G-, and other types of AFM stacking known for about 70 years.^[69] It represents alternating layers of FM-ordered chains parallel to each other within the plane but crossing and AFM-coupled in adjacent layers. Figure 1 illustrates the characteristic X-type AFM stacking with orthogonal A- and B-chains. Systems exhibiting symmetry between A- and B-chains under the condition $\mathbf{M}_{\text{A}} = -\mathbf{M}_{\text{B}}$ are classified as class-I X-type AFM stacking. This category corresponds to d-wave altermagnetism. Conversely, class-II X-type stacking

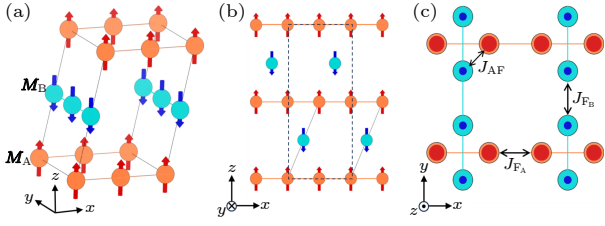


Fig. 1. Schematic of a crystal with X-type AFM stacking. (a) A primitive unit cell of the crystal. (b) Side view of the crystal, where the dashed box denotes the conventional unit cell and the solid box denotes the primitive unit cell. (c) Top view of two adjacent layers of the crystal.

arises from structural inequivalence between chains, manifested through distinct intrachain/interchain distances, or relative rotations/distortions of structural motifs. This category supports spin-splitting at the Γ point in compensated magnets.^[58,59,67] Note that there are also spin-degenerate antiferromagnets with class-III X-type AFM stacking where a magnetic unit cell doubles the crystal unit cell,^[67] which are not discussed in this work.

Figure 1(c) reveals a typical feature of the X-type AFM stacking, namely that the parallel FM chains are well isolated. As a result, the intralayer FM exchange coupling is dominated by the intrachain exchanges J_{FA} and J_{FB} , and the coupling between the parallel chains is negligible. The interlayer AFM exchange coupling J_{AF} connects the adjacent crossing chains and is usually weaker than J_{FA} and J_{FB} due to the interchain nearest-neighbor distance being larger than the intrachain distance.

We first discuss X-type antiferromagnets in the 2D limit with only two atomic layers exhibiting class-I X-type AFM stacking. Such a bilayer crystal represents an altermagnet with two FM-ordered chains of opposite magnetic sublattices connected by crystal symmetry. Using Hamiltonian (1), we calculate the magnon spectra as described in Supplementary Materials (SM). For class-I X-type AFM stacking, we consider the parameters $J_{FA} = J_{FB}$, $J_{FB} = J_F$, $J_{AF} = -\lambda J_F$, $K_A = \mu K_B$, and $K_B = 0.01 J_F$. In order to describe the slight inequivalence between chains for class-II X-type AFM stacking, we consider $J_{FA} = \mu J_{FB}$ with $\mu = 1.05$ as a representative. Figure 2(a) shows the results in the absence of the AFM coupling between the cross-chains ($\lambda = 0$). In this case, we find that the spin-split magnon band structure is similar to the electronic band structure.^[67] Specifically, the bands corresponding to the α -mode (β -mode) are strongly dispersive along the Γ -X (Γ -Y) direction but dispersionless along the Γ -Y (Γ -X) direction. This behavior reflects the fact that the α -mode (β -mode) is associated with the A-chain (B-chain) excitations and indicates that the magnon band structure is a superposition of the bands from the isolated A- and B-chains. As a result, the magnon density of states (DOS) shown in Fig. 2(d) (top panel) reveals van Hove singularities at band edges, revealing the 1D nature of the magnon bands.

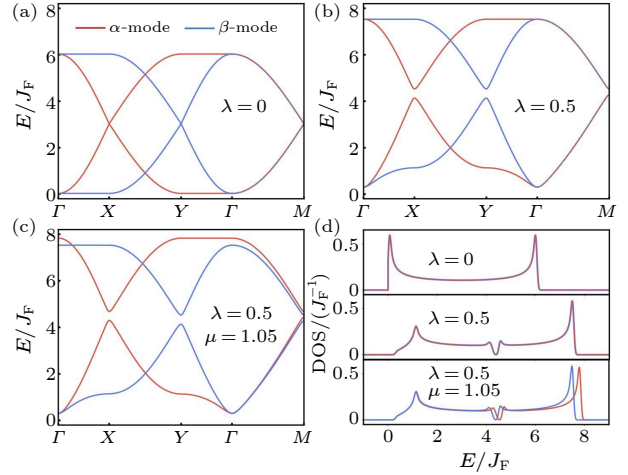


Fig. 2. The sublattice-resolved magnon band structure and DOS of an X-type AFM bilayer, for magnetic parameters $J_{FB} = J_F$, $J_{FA} = \mu J_{FB}$, $J_{AF} = -\lambda J_F$, $K_A = \mu K_B$, and $K_B = 0.01 J_F$. The magnon band structure of the AFM bilayer with class-I X-type stacking for (a) $\lambda = 0$, $\mu = 1$ and (b) $\lambda = 0.5$, $\mu = 1$. (c) The magnon band structure of the AFM bilayer with class-II X-type stacking for $\lambda = 0.5$, $\mu = 1.05$. (d) The DOS corresponding to the magnon band structure shown in (a)–(c).

As seen from Fig. 2(b), including an AFM interlayer exchange coupling in the calculation, $J_{AF} = -0.5 J_F$ ($\lambda = 0.5$), opens a band gap in the magnon spectra. This is accompanied by the bands becoming more (less) dispersive at low (high) energy. At the same time, we observe the increasing low-energy bandwidth due to a finite exchange field applied on both magnetic sublattices [Fig. 2(d), middle panel]. Nevertheless, the presence of the low-energy DOS peak even at sizable J_{AF} implies that the 1D nature of the magnon bands is still sustained.

For class-II X-type antiferromagnets (e.g., $J_{FA} = \mu J_{FB}$ with $\mu = 1.05$), we find a similar band structure [Fig. 2(c)], except that the degeneracy is removed at the Γ point,* which is a typical feature of a non-altermagnetic spin splitting.^[59] This leads to a non-degenerate sublattice-resolved magnon DOS with the peaks at the band edges remaining intact [Figs. 2(c) and 2(d)].

Such a strong chain-induced band anisotropy can produce a sizable thermal magnon spin current with a minimal ξ , provided that the temperature gradient is applied along one of the chain directions. To demonstrate this, we calculate thermal magnon spin conductivities σ_{nn}^s driven by the temperature gradient $\nabla_n T$ along the A- ($n = x$) or B- ($n = y$) chain. Figure 3(a) shows the calculated $\sigma_{xx}^s(\mathbf{k})$ projected onto the respective magnon band structures for $\lambda = 0.5$. Here we fix a constant $\tau_0 = 10$ ps widely used in literature^[60,70,71] for qualitative comparison across different magnetic systems. We note that the relaxation time of magnons τ_0 varies with momentum and temperature in real magnets. This should be carefully addressed in quantitative modeling of real materials.

It is seen that the thermal magnon spin conductivity is dominated by the bands originating from the chains

*The energy gap for the two bottom bands is small but non-vanishing, and is determined by the magnetocrystalline anisotropy difference between A- and B-chains, as shown in the SM.

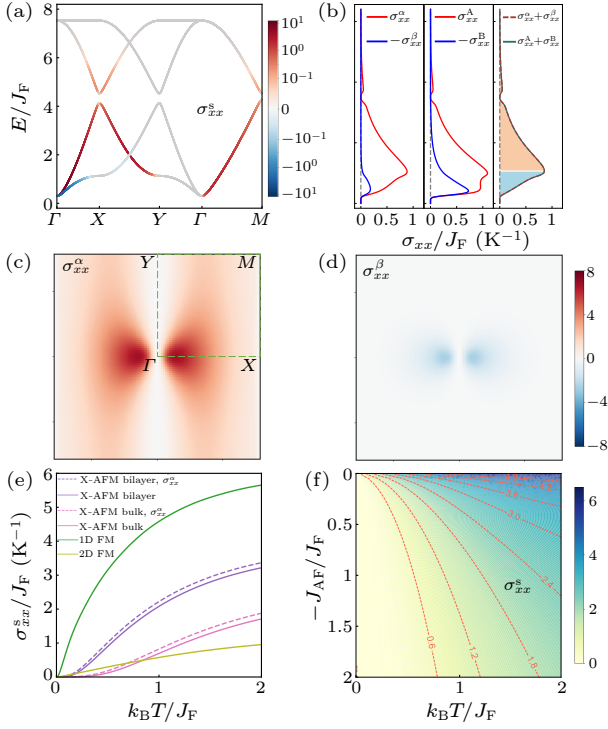


Fig. 3. (a) σ_{xx}^s projected magnon bands for an X-type AFM bilayer with $\lambda = 0.5$, $\mu = 1$, $k_B T = 0.8J_F$, and $\tau_0 = 10$ ps. (b) The energy distributions of the eigenmode-resolved magnon spin conductivity σ_{xx}^i (left), sublattice-resolved magnon spin conductivity σ_{xx}^α (middle), and the total spin conductivity σ_{xx}^s (right). The apricot and blue colors denote the energy windows with $|\sigma_{xx}^B/\sigma_{xx}^A|$ below and above 0.2, respectively. The $\mathbf{k}$ -dependent (c) $\sigma_{xx}^\alpha(\mathbf{k})$ and (d) $\sigma_{xx}^\beta(\mathbf{k})$ in the 2D Brillouin zone. (e) σ_{xx}^s as a function of temperature for the X-type AFM bulk and bilayer with $\lambda = 0.5$ and $\mu = 1$. The σ_{xx}^s for a 1D FM chain and a 2D FM square lattice are also included for comparison, where the nearest-neighbor FM exchange parameter is J_F . We also plot σ_{xx}^α of the X-type AFM bilayer as dashed lines for comparison. (f) A phase diagram of σ_{xx}^s with respect to J_{AF} and $k_B T$.

parallel to the temperature gradient. Specifically, as indicated by the large group velocity of α -mode along the Γ -X direction across the entire bandwidth, for $\nabla_n T$ along the x -direction, the α -mode controls $\sigma_{xx}^s(\mathbf{k})$, making it largely positive [red color in Fig. 3(a)]. While there is a nonvanishing contribution to $\sigma_{xx}^s(\mathbf{k})$ from β -mode in the less dispersive bottom band with finite group velocity in a narrow energy window at low energies [blue color in Fig. 3(a)], this contribution is small. This is confirmed by calculating the energy-resolved magnon spin conductivity $\sigma_{xx}^i(E) = \int \frac{d^2\mathbf{k}}{(2\pi)^2} \sigma_{xx}^i(\mathbf{k}) \delta[E - E(\mathbf{k})]$ where $\delta[E - E(\mathbf{k})]$ is a Dirac delta function. As shown in Fig. 3(b), $\sigma_{xx}^\beta(E)$ is finite only for the low energy, and $\sigma_{xx}^\alpha(E)$ is always dominant. These features are also seen from the distributions of $\sigma_{xx}^\alpha(\mathbf{k})$ and $\sigma_{xx}^\beta(\mathbf{k})$ in the 2D Brillouin zone [Figs. 3(c) and 3(d)], indicating that $\sigma_{xx}^\beta(\mathbf{k})$ is negligible compared to $\sigma_{xx}^\alpha(\mathbf{k})$. Notably, $\sigma_{xx}^\beta(\mathbf{k})$ vanishes along the Γ -Y and Γ -M directions despite the dispersive β -mode band along these directions [Figs. 2(b) and 3(a)-3(d)], clearly reflecting the unidirectional confinement induced by the 1D nature of X-type stacking. As a result, the compensation

of the thermal magnon spin current is largely suppressed by the X-type AFM stacking. This fact is also evident from Fig. 3(e), revealing a very small difference between σ_{xx}^s and σ_{xx}^α as a function of temperature [violet and magenta dashed and solid lines in Fig. 3(e)]. For example, at relatively high temperature $k_B T = 0.8J_F$, we obtain the magnon spin current compensation ratio of only $\xi = 8\%$. Considering the 1D nature of the magnon bands sustained with a sizable J_{AF} , one can conclude that the thermal magnon spin current is virtually controlled by a single magnetic sublattice, similar to that for a purely FM system.

To further verify the sublattice-dominant nature of magnon transport, we calculate the sublattice-resolved thermal spin conductivity σ_{nm}^A and σ_{nm}^B , which satisfy $\sigma_{nm}^s = \sigma_{nm}^\alpha + \sigma_{nm}^\beta = \sigma_{nm}^A + \sigma_{nm}^B$ (see the SM for details), as demonstrated in Fig. 3(b). We find that although σ_{xx}^B is non-negligible in the energy window corresponding to the less dispersive bottom bands along the Γ -X direction, it vanishes at higher energies. Consequently, $\sigma_{xx}^A \approx \sigma_{xx}^\alpha$ at these energies. This is highly nontrivial because, in an AFM-coupled system, spin wave modes on different sublattices are not eigenstates of the entire system; hence, σ_{nn}^i typically differs from σ_{nn}^α . Defining sublattice dominance as $|\sigma_{xx}^B/\sigma_{xx}^A| < 0.2$ and integrating $\sigma_{xx}^s(E)$ in the associated energy window [represented by the apricot area in the right panel of Fig. 3(b)], we observe that nearly 80% of the total σ_{xx}^s originates dominantly from the A-chain contribution. This result clearly demonstrates that the strong suppression of magnon spin current compensation is associated with the 1D nature of the chain-like sublattices inherent to the X-type stacking.

In addition, this sustained 1D nature further supports the enhanced magnon spin current due to its unidirectional confinement. We demonstrate this by comparing the calculated σ_{xx}^s to the magnon spin conductivities of a 1D FM chain and a 2D FM square lattice, where the nearest-neighbor FM exchange parameters are set to J_F . Here, in order to perform such comparisons of systems with different dimensionalities, we normalize the units of magnon spin conductivities presented in this work to those of a 2D system. Specifically, when calculating the conductivity of a 1D FM chain, we consider a 2D unit cell with the unit lattice constant identical to that of an X-type AFM bilayer, and the interchain exchange coupling is set to zero. For the bulk 3D X-type antiferromagnets, we multiply the calculated magnon spin conductivities by the interlayer distance. We find that although σ_{xx}^s of the X-type AFM bilayer is smaller than that of the 1D FM chain, it is much larger than that of the 2D FM square lattice [Fig. 3(e)]. This result clearly indicates that the 1D nature of X-type magnets supports giant magnon spin conductivity. Moreover, σ_{xx}^s remains notable even for a much larger J_{AF} , as reflected in the diagram of σ_{xx}^s with respect to J_{AF} and temperature [Fig. 3(f)]. This is also true for $\nabla_y T$ along the B-chain, where $\sigma_{yy}^s = -\sigma_{xx}^s$ (see the SM).

Going from a bilayer to a bulk X-type antiferromagnet with class-I AFM stacking, we find that σ_{xx}^s is somewhat reduced [Fig. 3(e) and the SM]. This is because each FM-ordered layer in the bulk has two adjacent layers (top and bottom) AFM-coupled to the former, while only one

is present in the bilayer. As a result, the effective AFM exchange field on each magnetic sublattice in the bulk is twice that in the bilayer. This enhances the contribution to σ_{xx}^s from the chains perpendicular to the temperature gradient. Nevertheless, we still observe weak compensation (e.g., $\xi = 20\%$ at $k_B T = 0.8 J_F$) and sizable σ_{xx}^s , which is larger than that of the 2D FM square lattice at not too low temperatures [Fig. 3(c)].

Similar magnon properties occur in class-II X-type antiferromagnets where A- and B-chains are asymmetric. In this case, σ_{xx}^s and σ_{yy}^s are slightly different, which remain sizable and uncompensated, as shown in Fig. S2 of the SM.

We further confirm the advantage of the X-type AFM stacking for magnon spin current generation by considering two uncompensated magnets with sizable net magnetizations. The first one is a class-I X-type AFM trilayer [Fig. 4(a)]. The second one is a bilayer ferrimagnet where the orthogonal A- and B-chains host non-equivalent moments $\mathbf{M}_A = -2\mathbf{M}_B = 2\mathbf{M}_0$, representing a class-II X-type AFM stacking [Fig. 4(b)]. Such systems can be prepared by layer-by-layer deposition using modern film-growth techniques.^[72,73] In both cases, the effective exchange fields on the sublattices are different from those in X-type antiferromagnets discussed above, which significantly influence the band structure (Fig. S4 of the SM). Despite the difference, the basic feature of the X-type AFM stacking is sustained, i.e., the bands dominated by A-chain (B-chain) excitations are strongly dispersive along the Γ -X (Γ -Y) direction but weakly dispersive along the Γ -Y (Γ -X) direction (Fig. S4 of the SM). As a result, σ_{xx}^s (σ_{yy}^s) is majorly contributed by the A-chains (B-chains).[†] As seen from Fig. 4, σ_{xx}^s is positive and much larger in these systems compared to the X-type AFM bilayer [Fig. 3(c)], due to the net magnetic moments on sublattice A being much larger than that on sublattice B. Notably, we find a negative σ_{yy}^s for both systems, indicating that a magnon spin current is polarized opposite to that expected by the net magnetic moment $\mathbf{M}_0$. This is in contrast to common expectation and serves as further evidence of the uncompensated magnon spin current nearly locked to the selected directions by the X-type AFM stacking.

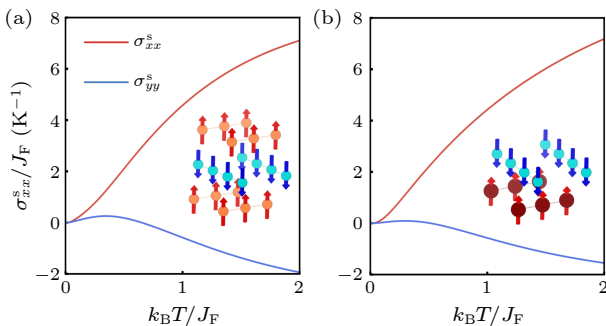


Fig. 4. (a) σ_{xx}^s and σ_{yy}^s as functions of temperature for an X-type AFM trilayer with $\lambda = 0.5, \mu = 1$. (b) σ_{xx}^s and σ_{yy}^s as functions of temperature for an X-type FIM bilayer with $\lambda = 0.5, \mu = 1$, and $\mathbf{M}_A = -2\mathbf{M}_B$.

[†]In a trilayer X-type antiferromagnet, there are six moments in a unit cell with four in the A-chain and two in the B-chain. In this case, four magnon eigenmodes are excited mainly by the A-chain, and two magnon eigenmodes are excited mainly by the B-chain. We independently calculate the contribution of each eigenmode to thermal spin transport and sum up these contributions to obtain the magnon spin conductivities.

When the temperature gradient $\nabla_m T$ is along direction m , making an angle θ relative to the x -axis, a transverse magnon spin current in X-type magnets appears along the direction $n \perp m$ with conductivity $\sigma_{nm}^s = (\sigma_{xx}^s - \sigma_{yy}^s) \sin \theta \cos \theta$, representing a spin Nernst effect. The maximum $|\sigma_{nm}^s| = |(\sigma_{xx}^s - \sigma_{yy}^s)/2|$ occurs for m along the diagonal directions (Fig. S5 of the SM). Since σ_{xx}^s and σ_{yy}^s have opposite signs and are both large, we expect that the transverse magnon spin current driven by the X-type AFM stacking will be much stronger than that originating from the Berry curvature.

The giant uncompensated magnon spin currents in X-type magnets are robust even when spin-orbit coupling (SOC) is included. This is demonstrated by the slight variation of the magnon spin conductivity calculated with consideration of the exaggerated Dzyaloshinskii-Moriya interaction (DMI) parameters and different magnetic anisotropies, as shown in Fig. S7 of the SM.

We emphasize that the key property of the X-type AFM stacking resulting in the predicted giant uncompensated magnon spin current is the strong intra-sublattice FM exchange coupling in one direction and negligible coupling in the perpendicular direction. This property is a direct consequence of the chain-like sublattice structure of the X-type AFM stacking producing strong anisotropy between the two magnetic sublattices. The sublattice anisotropy has implications for understanding the thermal magnon spin current generation in AFM-stacked magnets. For example, rutile altermagnets do not exhibit sufficiently strong in-plane structural anisotropy within their magnetic sublattice, resulting in the strong compensation of the magnon spin current (e.g., $\xi = 87\%$ for RuO_2 at 300 K, see the SM).^[61,62] In other altermagnetic candidates, such as M_2X_2O ($M = \text{V, Cr}; X = \text{Se, Te}$),^[43,60] the intra-sublattice nearest neighbors are connected by M -X- M bonds in one direction and M -O- M bonds in perpendicular directions. This sublattice anisotropy makes the FM exchange constants along these directions very different, supporting a not-too-large ξ and thus a sizable magnon spin current (see the SM).^[60] The X-type magnets can be regarded as an extreme case of the AFM-stacked magnets with the strongest structural anisotropy of magnetic sublattices. For example, in the case of the typical X-type antiferromagnet $\beta\text{-Fe}_2\text{PO}_5$,^[67,74,75] which is described by the magnetic parameters $J_{FA} = J_{FB} = 32 \text{ meV}$, $J_{AF} = -9.2 \text{ meV}$, and $K_A = K_B = -0.44 \text{ meV}$, we calculate tiny $\xi = 6\%$ and giant $\sigma_{xx}^s = -\sigma_{yy}^s = 30.8 \text{ meV/K}$ at 300 K (details in the SM). This implies an outstanding efficiency of the magnon spin current generation of this X-type antiferromagnet compared to the known AFM-stacked magnets.

Overall, X-type magnets provide conditions for the efficient magnon spin current generation, which is usually difficult in antiferromagnets and ferrimagnets. Due to the unique orthogonal cross-chain structure, the counteraction of the two magnetic sublattices in these magnets is effectively suppressed, supporting magnon spin currents with

minimal compensation. The 1D nature of the isolated FM chains significantly enhances the magnon spin conductivity and, in combination with the minimal sublattice compensation, makes it giant. These unprecedented properties of X-type magnets are promising for applications in low-power high-performance spintronics.

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