

***D*-Wave Phonon Angular Momentum Texture in Altermagnets by Magnon-Phonon Hybridization**

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In altermagnets, the magnon bands are anisotropically spin split in reciprocal space without relativistic or dipolar spin-spin interactions. In this Letter, we theoretically study magnons and phonons coupled by spin-lattice interaction in a two-dimensional square-lattice *d*-wave altermagnet. We show that phonon-chirality-selective magnon-phonon hybridization can be caused by interfacial Dzyaloshinskii-Moriya interaction leading to the emergence of hybrid quasiparticles that possess finite phonon angular momentum. These hybrid quasiparticles are called magnon polarons and consist of spin-polarized magnons and chiral phonons. Their phonon angular momentum texture follows the *d*-wave character of the magnon spin texture, opening up the possibility of phononic counterparts to the electronic response effects in altermagnets, such as a *phonon angular momentum splitter effect*, i.e., the generation of a transverse phonon angular momentum current induced by a temperature gradient—the bosonic analog of the spin-splitter effect.

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Introduction—Chiral phonons [1–3], the quasiparticles of circularly polarized lattice vibrations, have recently been investigated due to a range of emerging phenomena. They have been found to contribute to spin relaxation [4,5], ultrafast demagnetization [6,7] and the Einstein–de Haas effect [8–10]. Notably, chiral phonons carry nonzero angular momentum and can therefore function as angular momentum carriers in nonmagnetic materials [11,12]. Here we are using the term “chiral phonon” synonymously to “axial phonon” proposed in Ref. [13] for phonons that carry angular momentum.

The systems in which chiral phonons occur still require extensive research. Chiral phonons may, for example, be found in lattices with broken inversion symmetry [1,2,14–16]. Due to the preserved time-reversal symmetry, the phonon angular momentum has an odd parity: $\mathbf{L}(\mathbf{k}) = -\mathbf{L}(-\mathbf{k})$ [15]. In order to realize *even-parity* textures, time-reversal symmetry must be broken [15], e.g., by the coupling to magnons, the quasiparticles of spin excitations [17–23].

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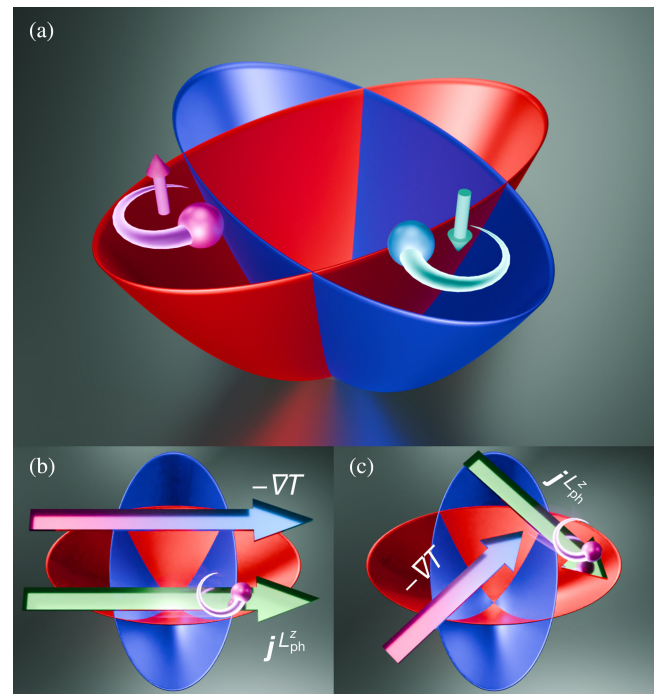


FIG. 1. (a) Generation of phonon angular momentum in altermagnets through phonon coupling with the underlying spin-split magnon band structure. Their hybridization selectively depends on the circular polarization of the phonons and, thus, it separates left- and right-handed phonons imprinting an alternating phonon angular momentum texture in reciprocal space. (b),(c) Visualization of the phonon angular momentum splitter effect, i.e., the generation of a phonon angular momentum current by the application of a temperature gradient.

In this Letter, we delineate an approach to design even-parity phonon angular momentum textures beyond s wave. Our approach is based on a chirality-selective magnon-phonon coupling, which leads to the hybridization between magnons and phonons. In particular, we show that the strength of magnon-phonon coupling originating from interfacial Dzyaloshinskii-Moriya interaction (DMI) depends on the magnon spin and the phonon angular momentum. We study their hybridization in the context of a d -wave altermagnet [24,25] and demonstrate how the symmetry of the magnon spin texture can be imprinted onto the phonon angular momentum texture of the emergent magnon polarons [cf. Fig. 1(a)]. Our findings encourage the exploration of the vibrational degrees of freedom in novel unconventional magnetic phases such as altermagnets or antialtermagnets [26] and highlight the potential of unconventional magnetism for realizing chiral phonons.

Results—We consider the minimal, two-dimensional altermagnet depicted in Fig. 2(a) with the two sublattices $\uparrow$ (red circles) and $\downarrow$ (blue circles). The localized spins assume a collinear Néel order and interact via antiferromagnetic nearest and anisotropic ferromagnetic next-nearest neighbor Heisenberg coupling (cf. End Matter). Using linear spin-wave theory [27], we obtain the dispersion relation depicted as red and blue lines in Fig. 2(b), which features spin-split magnon modes. Their splitting is maximal and opposite along the high-symmetry lines $\overline{\Gamma X}$ and $\overline{\Gamma Y}$ and vanishes along the nodal lines $\overline{\Gamma M}$. Thus, the model exhibits the characteristic features in the magnon band structure expected in altermagnets [23,34,35].

The phonon Hamiltonian of the system respects the symmetry of the square lattice and comprises nearest-neighbor interactions with restoring forces along longitudinal $K_L^{(1)}$ and transverse K_T directions with respect to the bonds as well as longitudinal next-nearest neighbor interactions $K_L^{(2)}$ (cf. End Matter). The phonon dispersion relation is depicted in gray in Fig. 2(b) along with the magnon modes. Note that in the structural unit cell the phonon system comprises only a single atom per unit cell forming two acoustic bands. However, since the magnetic unit cell is larger, the corresponding Brillouin zone is smaller and the bands have to be down-folded yielding four phonon branches. In general, the checkerboard geometry invoked for the magnons also modifies the phonons, but it does not qualitatively change the results presented here [27]. Because of the restriction to in-plane displacements, the phonon angular momentum is aligned with the z axis and can be expressed as [8]

$$L_{\text{ph}}^z = \sum_i (u_i^x p_i^y - u_i^y p_i^x). \quad (1)$$

Symmetry demands that the expectation values of L_{ph}^z vanish in an inversion- and time-reversal-symmetric (non-magnetic) square lattice [15].

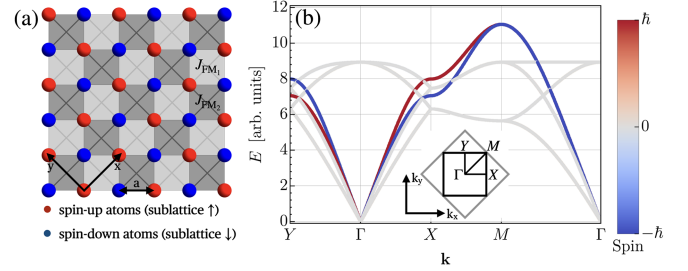


FIG. 2. (a) Altermagnetic model comprising antiferromagnetic nearest-neighbor interactions and directional, ferromagnetic next-nearest neighbor interactions as indicated by the gray checker-board pattern. (b) Magnon and phonon modes in the absence of spin-lattice coupling. The color indicates the magnon spin. Phonon modes are depicted in gray while the red and blue colors confirm the spin-dependent splitting of the magnon modes. Inset: Brillouin zone of the magnetic (structural) unit cell shown in black (gray). (The coordinate systems of the real and reciprocal spaces are rotated by 45° to each other.) The parameters used for the calculations are $J_{\text{AFM}} = -1$, $J_{\text{FM}_1} = 0.383$, $J_{\text{FM}_2} = 0.5$, $S = 1$, $M = 10$, $K_L^{(1)} = 160$, $K_T = K_L^{(2)} = 40$, and $a = 1$.

Turning to the magnon-phonon coupling, the nonrelativistic altermagnet does not exhibit magnon-phonon hybridization because exchange magnetostriiction preserves the magnon spin. To break spin conservation, we additionally invoke DMI, which is the leading-order effect of spin-orbit interaction [36,37]. In particular, we study interfacial DMI $\hat{H}_{\text{DMI}} \propto \mathbf{D}_{ij} \cdot (\mathbf{S}_i \times \mathbf{S}_j)$, which results from the breaking of the out-of-plane mirror symmetry. Since the orientation of the DM vector $\mathbf{D}_{ij} = D\hat{\mathbf{R}}_{ij} \times \hat{\mathbf{z}}$ is in the plane and perpendicular to the bond vector $\mathbf{R}_{ij}$, it is also perpendicular to the Néel vector and does not modify the noninteracting magnon dispersion relation [38]. (Although the true ground state in the presence of DMI would be a noncollinear state [39], an easy-axis anisotropy would stabilize the collinear ground state for sufficiently small D . We do not implement the anisotropy for simplicity, since it does not change the qualitative effects discussed in the following [27].) The influence of the lattice excitations on the magnons enters the model by the modification of the bond vectors $\mathbf{R}_{ij} \mapsto \mathbf{R}_{ij} + \mathbf{u}_{ij}$. Expansion of the DMI up to first order in the displacements of the atoms yields the coupling Hamiltonian [17,19,27,38]

$$\hat{H}_{\text{mpc}} = \frac{DS}{R} \sum_{\langle\langle i,j \rangle\rangle} \{\pm(\mathbf{u}_i - \mathbf{u}_j) \cdot (\mathbf{S}_i - \mathbf{S}_j)\}. \quad (2)$$

The sum in the expression of $\hat{H}_{\text{mpc}}$ runs over all next-nearest neighbors $\langle\langle i,j \rangle\rangle$. The “+” sign holds for those sites i, j of the $\uparrow$ and the “−” sign for the $\downarrow$ sublattice. We note in passing that we have neglected contributions that originate from the distance dependence of $D(R)$ [38], which we explicitly analyze in Supplemental Material [27].

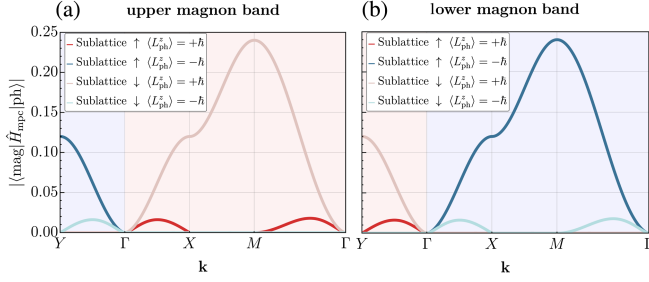


FIG. 3. Magnon-phonon coupling strength between the (a) upper, (b) lower magnon branch and sublattice-polarized phonon modes. All lines represent the absolute value of one of the 8 matrix elements of $\hat{H}_{\text{mpc}}$ [Eq. (2)] represented in the basis of the magnon bands and the sublattice-resolved chiral phonons (cf. Supplemental Material [27]). The background color and the line color indicate the signs of the magnon spin and the phonon angular momentum, respectively. The matching of colors indicates the chiral selectivity of the magnon-phonon coupling $\hat{H}_{\text{mpc}}$.

To characterize $\hat{H}_{\text{mpc}}$, we project it onto the magnon bands as eigenmodes of the bare magnon Hamiltonian (lower magnon band, upper magnon band) and sublattice-localized chiral phonons (right- and left-handed chiral phonons localized on sublattices $\uparrow$ and $\downarrow$) and take the absolute value of the matrix elements representing all eight interaction channels. The mathematical details of this analysis are described in Supplemental Material [27] and the result is plotted in Fig. 3. The background and the line colors indicate the sign of the magnon spin and phonon angular momentum, respectively. For the upper magnon band [Fig. 3(a)], the magnon spin is negative along $\overline{\Gamma Y}$ (indicated by blue background color), and the band couples exclusively to left-handed circularly polarized phonons with negative angular momentum ($\langle L_{\text{ph}}^z \rangle = -\hbar$; indicated by blue line color) on both sublattices. The coupling strength on the $\downarrow$ sublattice (faint color) is weaker than on the $\uparrow$ sublattice (vivid color) and vanishes at the Y point. Along $\overline{\Gamma X}$ the magnon spin reverses sign (red background color), such that the band couples solely to right-handed circularly polarized phonons with positive angular momentum ($\langle L_{\text{ph}}^z \rangle = \hbar$; red line color). Along $\overline{\Gamma M}$, the two magnon branches become degenerate, rendering their labeling arbitrary. Overall, $\hat{H}_{\text{mpc}}$ couples each magnon spin channel selectively to a single phonon circular polarization, determined by the sign of the magnon spin.

The physical origin of this selectivity can be intuitively understood from Eq. (2). In antiferromagnets and altermagnets the two magnon branches have opposite chiralities, corresponding to clockwise or counter-clockwise spin precession in the xy plane [34,40]. The dynamic spin precession induces an effective field acting on the lattice displacements through $\hat{H}_{\text{mpc}}$. For corotating chiral phonons—those exhibiting the same sense of rotation as the spin precession—the time average of Eq. (2) remains

finite leading to an energy shift. In contrast, the coupling for counter-rotating phonons—those exhibiting the opposite sense of rotation as the spin precession—averages to zero explaining intuitively the chiral selection rule of the coupling. Linearly or elliptically polarized phonons should be regarded as superpositions of left- and right-handed circular polarizations, which therefore couple to both magnon bands and the coupling strength depends on the proportion of the compatible circular polarization. Only for chiral phonon modes featuring the incompatible circular polarization, the coupling is zero. Within a given magnon band, the sublattice amplitudes of the spin precession are generally unequal. For instance, the spin-down magnon mode exhibits a larger precession amplitude on the $\uparrow$ sublattice, as the magnon spin is defined as the difference of the spin expectation values between the excited and the ground states [41]. As $\mathbf{k}$ approaches the boundary of the first Brillouin zone, the amplitude of the $\downarrow$ sublattice diminishes, while that of the $\uparrow$ sublattice remains finite.

The dispersion relation of the coupled system

$$\hat{H} = \hat{H}_{\text{m}} + \hat{H}_{\text{ph}} + \hat{H}_{\text{mpc}} \quad (3)$$

is plotted in Fig. 4(a). Here, the color indicates the magnon/phonon/mixed character of the quasiparticles. Evidently, avoided crossings emerge at the intersection of magnon and phonon bands, giving rise to hybrid magnetic and structural excitations dubbed “magnon polarons” [42]. Furthermore, the close-up shown in Fig. 4(b) into the highlighted section as part of the $\overline{\Gamma Y}$ line provides insights into the generation of phonon angular momentum by magnon-phonon coupling through time-reversal symmetry breaking. In agreement with previous studies, a distinct avoided crossing pattern emerges at the intersection of two phonon bands and a single magnon band [19,43]. While the energy and phonon character of the central mode remain unchanged, two magnon-polaron modes emerge whose energies are shifted. This is a general feature of three-band magnon-phonon hybridization that one state is uncoupled to the magnon and remains a pure phonon, while the other two states are magnon polarons with equal proportions of magnons and phonons [27]. Additionally, the hybridization generates phonon angular momentum for the states close to the avoided crossing [cf. Fig. 4(c)]. Interestingly, we find that the phonon angular momentum changes sign and vanishes directly at the avoided crossing. Only in its vicinity it becomes nonzero. Further analysis reveals that the sublattice-resolved phonon angular momentum is maximized directly at the avoided crossing reaching values of up to $\hbar/2$ [44] and the lattice vibrations are almost fully circularly polarized [cf. Figs. 4(e) and 4(f)]. Because the circular polarization is reversed on the two sublattices, the sublattice-resolved angular momenta partially compensate each other and the net angular momentum is strongly suppressed.

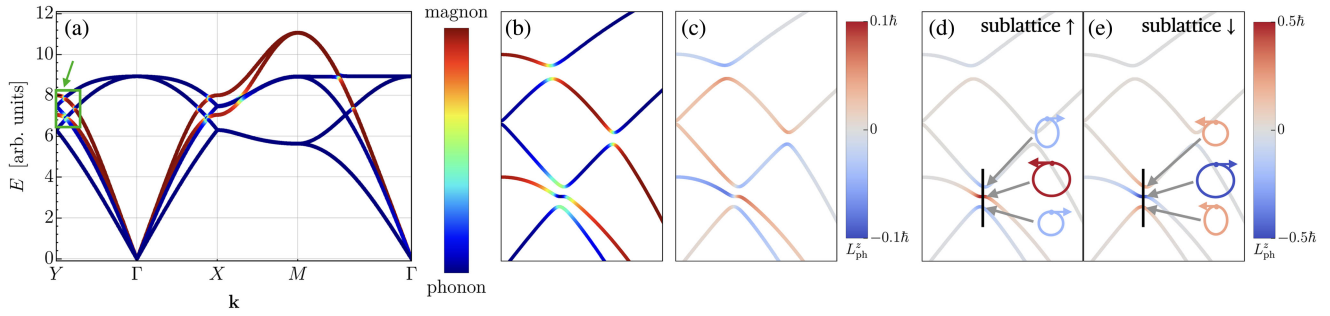


FIG. 4. Coupled magnon-phonon hybrid system. (a) Band structure of the magnon polarons along a high-symmetry path in the magnetic Brillouin zone. The color indicates the magnon/phonon character of the quasiparticles. The region highlighted by the green box and arrow is displayed in greater detail in (b)–(e). (b) Zoom of the avoided crossing pattern and quasiparticle character indicated by color. (c) Illustration of the same region but depicting the phonon angular momentum. (d),(e) Sublattice-resolved phonon angular momentum for sublattice $\uparrow$ and sublattice $\downarrow$, respectively. The insets demonstrate the real-space vibrational motion (amplitude and phase) of the respective sublattices. The parameters used for the calculations are $J_{\text{AFM}} = -1$, $J_{\text{FM}_1} = 0.383$, $J_{\text{FM}_2} = 0.5$, $S = 1$, $M = 10$, $K_L^{(1)} = 160$, $K_T = K_L^{(2)} = 40$, $D = 0.25$, and $a = 1$.

To analyze the cause of the compensation of the phonon angular momentum, we focus on the crossing of the lower magnon band with two phonon bands in Fig. 4(b). Figure 3 indicates that the lower magnon mode couples almost exclusively to phonons on the down-sublattice ($\downarrow$) and with a phonon angular momentum of $+\hbar$. Figure 4(e) displays the phonon angular momentum of the $\downarrow$ sublattice. Indeed, the magnon polaron bands exhibit a positive phonon angular momentum, while the uncoupled phonon mode has a negative phonon angular momentum. At the (avoided) crossing the chirality-selective magnon-phonon coupling superimposes the linearly polarized phonons to right- and left-handed chiral phonons and splits them by selectively hybridizing the magnon and the phonon with matching handedness. The uncoupled phonon consequently has opposite handedness. On the other hand, the opposite phonon angular momentum on the $\uparrow$ sublattice can be explained by the phase differences between the sublattices of the crossing phonon bands (cf. End Matter).

Finally, Fig. 5 depicts the total phonon angular momentum of all bands across the entire Brillouin zone. The d -wave structure of the phonon angular momentum texture can be clearly observed. The sign of the angular momentum alternates in the Brillouin zone and the phonon angular momentum vanishes along the two nodal lines at $\overline{\Gamma M}$. One can distinguish two kinds of distributions: there are broad, diffuse features as well as sharp line-shaped “hotspots.” The former can be mostly recognized in areas where the band either has magnon character (e.g., inner part of band 4) or where the band is close to another magnon band (e.g., outer area of band 2). The pronounced line-shaped features originate from avoided crossings between magnon and phonon bands. Since the energy gaps are relatively small, the lines have small thicknesses. For regions where the band exhibits phonon character and is energetically separated from magnon bands, the phonon angular momentum

is vanishingly small (e.g., band 6 or inner areas of band 1 and 2).

Discussion and conclusion—The alternating sign of the phonon angular momentum across the Brillouin zone represents a qualitatively new form of chiral phonon behavior, distinct from previously studied uniform or valley-contrasting chiral phonons [1,8,15]. This momentum-dependent angular-momentum texture directly results from the interplay between the underlying d -wave magnetic order and the chiral selection rules of the relativistic magnon-phonon coupling. In our analysis we have focused on a particular chirality-selective part of the interfacial DMI, but chirality-selective forms of magnon-phonon coupling can emerge more generally in the framework of rotationally invariant spin-lattice coupling [45–47]. The considered interfacial DMI may generate additional terms stemming from the distance dependence of $D(R)$ lifting the perfect chiral selectivity in a k -dependent fashion. While it renders the phonon angular momentum texture more complex, its general features remain robust as we demonstrate in Supplemental Material [27]. The discovery of d -wave chiral phonons lays the foundation for realizing phononic counterparts to electronic responses in altermagnets.

d -wave altermagnets are known to exhibit a spin-splitter effect [48]. Based on the symmetry we can predict a *phonon angular momentum splitter effect*, where a temperature gradient generates a flow of phonon angular momentum. If the temperature gradient is applied along a high-symmetry spin-polarized direction, the phonon angular momentum current is parallel or antiparallel to the gradient [Fig. 1(b)]. If the temperature gradient is applied along a nodal line, the phonon angular momentum current is transversal to the gradient [Fig. 1(c)]. This nontrivial angular dependence distinguishes the splitter effect from the phonon angular momentum Hall effect proposed by

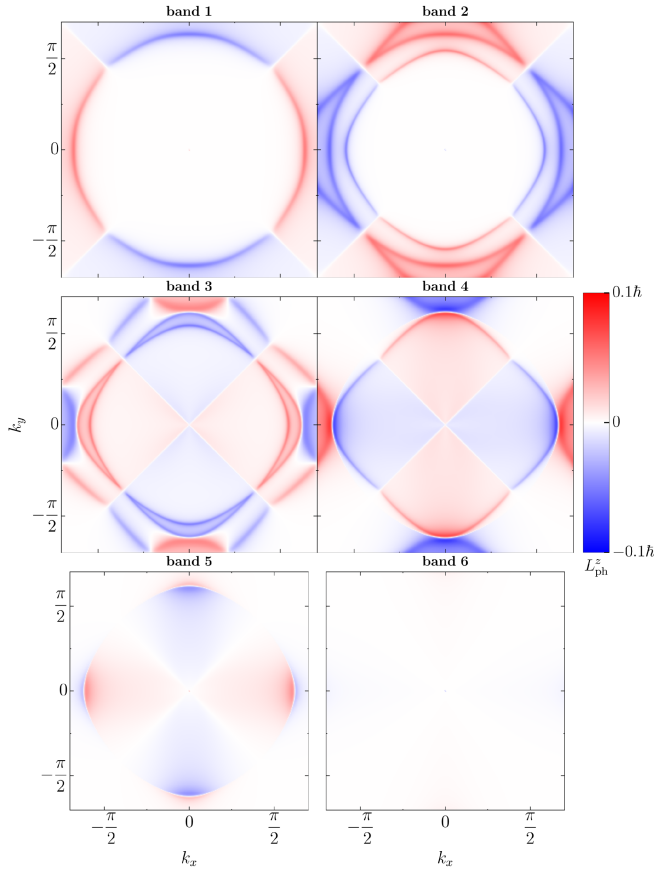


FIG. 5. Phonon angular momentum for all bands across the entire Brillouin zone. Nonzero angular momentum occurs around regions of avoided crossings and clearly reflects the d -wave character of the altermagnet. Bands are labeled in descending order starting with the highest energy. Parameters same as in Fig. 4.

Park and Yang [11]. Alternatively, the behavior under time reversal enables experimental differentiation of the two because the Hall effect remains unchanged and the splitter effect changes its sign upon reversal of the Néel order. Furthermore, the splitter effect is proportional to the relaxation time and dominates in clean samples, whereas dirty samples favor the intrinsic Hall effect [49]. Experimental studies have reported on measurements of phonon angular momentum currents via voltage generation in heavy-metal contacts [50], delineating a potential route for observing the phonon angular momentum splitter effect.

The applications are not limited to the splitter effect. Altermagnets form a net magnetization when subjected to strain, which is known as piezomagnetism [51,52]. Since in ionic crystals the phonon angular momentum is accompanied by an orbital magnetic moment [53,54], phonon piezomagnetism contributes in altermagnets. In addition, phononic analogs of spin filters and spin torques [55,56], and the nonlinear thermal Edelstein effect [57,58] enrich the research landscape for (d -wave) chiral phonons.

A general challenge for measuring phononic analogs of electronic and magnonic responses in altermagnets lies in their apparent symmetry equivalence. For instance, observing the phonon angular momentum splitter effect requires disentangling it from the concomitant magnon spin splitter effect [23,59]. Crucially, however, the two phenomena originate from distinct mechanisms: the magnon spin texture is nonrelativistic, whereas the phonon angular momentum is inherently relativistic. Consequently, the magnon spin texture is constrained by the spin point group (neglecting relativistic corrections), while the phonon angular momentum is fundamentally governed by the magnetic point group. This implies that, in contrast to the robust magnon spin, the phonon angular momentum depends sensitively on the Néel vector orientation. A sufficiently strong magnetic field can reorient the Néel vector away from its easy axis, thereby changing the plane of spin precession. This, in turn, alters which spin components couple to the phonons and thus controls the resulting phonon angular momentum. Within the model presented here, an in-plane Néel vector eliminates the phonon angular momentum because out-of-plane spin precession cannot generate out-of-plane atomic motion in a two-dimensional system. A weaker magnetic field applied along the easy axis shifts the magnon bands and therefore generates s -wave components for the chiral phonons. Hence, chiral phonons generated by magnetism offer great tunability compared to other approaches. In addition, the sensitivity of the effect to interfacial DMI suggests possible routes for electrically or structurally tunable phononic devices, which may be relevant for phononics and hybrid spin–lattice information processing. This highlights altermagnets as tunable platforms for realizing chiral phonons.

Our theory is valid for insulating altermagnets with sufficient relativistic magnetoelastic coupling. Among two-dimensional d -wave materials, monolayer $\text{Cr}_2\text{Te}_2\text{O}$ and $\text{Cr}_2\text{Se}_2\text{O}$ [60] realize the symmetry of our spin Hamiltonian (space group $P4/mmm$), but rely on a substrate that induces interfacial DMI. Another possibility are Janus altermagnets such as monolayer V_2SeTeO [61]. Here, the substitution of Se by Te in one layer breaks the out-of-plane mirror symmetry in the magnetic V layer and induces interfacial DMI intrinsically. Our results can also be generalized to three dimensions, where many more candidate materials have been identified [62]. Although we have focused on a d -wave altermagnet, imprinting phonon angular momentum by magnon-phonon coupling with chiral selection rules can be applied to g - and i -wave magnets as well, where higher-wave phonon angular momentum textures are generated.

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Data availability—Upon reasonable request, the data associated with this article can be made available on Zenodo [63].

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End Matter

Magnon Hamiltonian—The spin Hamiltonian reads

$$\hat{H}_m = -\frac{J_{\text{AFM}}}{2} \sum_{\langle i,j \rangle} \mathbf{S}_i \cdot \mathbf{S}_j - \frac{J_{\text{FM}_1}}{2} \sum_{\langle\langle i,j \rangle\rangle_1} \mathbf{S}_i \cdot \mathbf{S}_j - \frac{J_{\text{FM}_2}}{2} \sum_{\langle\langle i,j \rangle\rangle_2} \mathbf{S}_i \cdot \mathbf{S}_j. \quad (\text{A1})$$

Here, $J_{\text{AFM}} < 0$ and $J_{\text{FM}_{1,2}} > 0$ describe the antiferromagnetic and ferromagnetic interaction, respectively. $\langle i, j \rangle$ runs over all nearest neighbors, while $\langle\langle i, j \rangle\rangle_l$ runs over all next-nearest neighbors with bonds located within light gray ($l = 1$) or dark gray squares ($l = 2$) in Fig. 2(a). The splitting of the magnon bands is given by

$$\Delta E = 2\hbar^2 S (J_{\text{FM}_2} - J_{\text{FM}_1}) [\cos(\sqrt{2}k_y a) - \cos(\sqrt{2}k_x a)]. \quad (\text{A2})$$

Phonon Hamiltonian—The dynamics of the atoms is governed by

$$\begin{aligned} \hat{H}_{\text{ph}} = & \sum_i \frac{\mathbf{p}_i^2}{2M} + \frac{1}{2} \sum_{\langle i,j \rangle} \left[(K_{\text{L}}^{(1)} - K_{\text{T}}) (\mathbf{u}_{ij} \cdot \hat{\mathbf{R}}_{ij})^2 + K_{\text{T}} \mathbf{u}_{ij}^2 \right] \\ & + \frac{K_{\text{L}}^{(2)}}{2} \sum_{\langle\langle i,j \rangle\rangle} (\mathbf{u}_{ij} \cdot \hat{\mathbf{R}}_{ij})^2, \end{aligned} \quad (\text{B1})$$

where $\mathbf{u}_{ij} = \mathbf{u}_j - \mathbf{u}_i$ and $\mathbf{p}_i$ and $\mathbf{u}_i$ are the momentum and position operator of an atom at lattice site i , respectively, describing the dynamics in the xy plane. $\hat{\mathbf{R}}_{ij}$ refers to the unit bond vector between atoms at lattice sites i and j , while M is the mass of the nuclei.

Compensation of the phonon angular momentum—Due to the magnetism-induced downfolding, the longitudinal acoustic crosses with the transversal optical phonon branch, for which the atoms in the unit cell oscillate in phase and out of phase, respectively. Mathematically, the phases are given by $\exp\{i(\mathbf{k} + \mathbf{G}) \cdot \Delta \mathbf{b}\}$ where $\Delta \mathbf{b} = [a/\sqrt{2}](110)^{\text{T}}$ is the bond vector between the basis atoms and $\mathbf{G}$ is either zero for the acoustic branch, or, for an optical branch, an appropriate reciprocal lattice vector of the magnetic Bravais lattice. By choosing an arbitrary $\mathbf{k}$ point along $\overline{\Gamma\text{Y}}$, represented by $\mathbf{k} = (0 \ k_y \ 0)^{\text{T}}$ and by setting $\mathbf{G} = [0 \ (2\pi/\sqrt{2}a) \ 0]^{\text{T}}$, one finds $\mathbf{u}_{\downarrow} = \pm e^{-i\Phi} \mathbf{u}_{\uparrow}$ for the acoustic and optical branches, where $\mathbf{u}_p$ is the wave function of basis atoms $p = \uparrow, \downarrow$ and $\Phi = -(a/\sqrt{2})k_y$. The longitudinal $|l\rangle = (0 \ 1 \ 0 \ e^{i\Phi})^{\text{T}}$ and transverse modes $|t\rangle = (1 \ 0 \ -e^{i\Phi} \ 0)^{\text{T}}$ are linearly polarized along y and x direction, respectively. Here, we represent the wave functions in the $(u_{\downarrow}^x \ u_{\downarrow}^y \ u_{\uparrow}^x \ u_{\uparrow}^y)$ basis. If we now construct right-handed circularly polarized phonons ($\langle L_{\text{ph}}^z \rangle = +\hbar$) on the $\downarrow$ sublattice, the $\uparrow$ sublattice will be left-handed circularly polarized:

$$|t\rangle + i|l\rangle = (1 \ i \ -e^{i\Phi} \ ie^{i\Phi})^{\text{T}}, \quad (\text{C1})$$

and vice versa. Hence, by selectively coupling a certain circularly polarized phonon mode of the $\downarrow$ sublattice to the magnon band, the opposite chirality is naturally imposed on the $\uparrow$ sublattice. Since the magnon, depending on $\mathbf{k}$, couples to both sublattices, the phonons are not circular, but elliptical. The closer the crossing is located at the edge of the Brillouin zone, the more the magnons specifically couple to one sublattice (in this case $\downarrow$) and the ellipticity is modified.