High-resolution polarised Raman scattering study on spin-phonon coupling in multiferroic MnWO₄

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Abstract
MnWO₄ passes three antiferromagnetic phase transitions below the Néel temperature \( T_\text{N} = 13.5 \) K. One of these phases is also ferroelectric and MnWO₄ is therefore referred to as multiferroic. We investigated the effect of magnetic and electric ordering on the lattice dynamics by high-resolution polarized Raman spectroscopy. With an accuracy of about 0.02 cm\(^{-1}\), anomalous shifts of nearly all Raman-active phonon wavenumbers were observed below \( T_\text{N} \). High-energy metal–oxygen stretching modes soften on cooling, while low-energy vibrations become harder. Due to their linear temperature dependence, these anomalies can be attributed to a higher-order spin–phonon coupling process. The phonon linewidths are unaffected by the magnetic phase transitions and continuously increase with temperature as expected for three-phonon decay processes.

1. Introduction

Multiferroic materials have gained attention during recent years, since they provide the possibility of magnetic and electric ordering in a single phase. In particular, multiferroics with a cycloidal spin arrangement are of interest because the spontaneous polarization is directly connected to the magnetic order [1–3]. One example for this type of compound is MnWO₄, which is monoclinic and crystallizes in the space group \( P2_1/c \) [4] with the lattice constants \( a = 4.8226(3) \) Å, \( b = 5.7533(6) \) Å, \( c = 4.9923(6) \) Å, \( \beta = 91.075(7)^\circ \) at 1.5 K [5]. At high temperatures it is paramagnetic and paraelectric and passes successively three magnetic phase transitions on cooling [5–7]. At \( T_\text{N} \approx 13.5 \) K, an incommensurate antiferromagnetically ordered AF3 phase is entered, with a slightly temperature dependent propagation vector \( \mathbf{q} = (-0.2145, 0.5, 0.4580) \) [8, 9]. The magnetic moments are sinusoidally modulated along the easy axis within the \( ac \)-plane at an angle of 35° to the \( a \)-axis. At \( T_2 \approx 12.5 \) K another antiferromagnetic phase, AF2, is entered, which is characterized by an additional \( b \)-component of the magnetic moments, giving rise to a helical spin arrangement with a cycloidal component and the same propagation vector as in AF3.

While the first two phase transitions are of second order, the first order lock-in transition at \( T_1 \approx 7 \) K leads to the phase AF1. Here, the magnetic moments are again aligned along the easy-axis without any \( b \)-component. The propagation vector locks into the commensurate value of \( \mathbf{q} = (-0.25, 0.5, 0.5) \), corresponding to a 16-fold magnetic superstructure [5]. Only the intermediate AF2 phase is multiferroic due to the interaction between the non-collinear spin arrangement and the oxygen ions. The inverse Dzyaloshinski–Moriya interaction [2, 10, 11] leads to a displacement of the oxygen ions and an electric polarization along the \( b \)-direction. By applying an electric field, the polarization and equally the chirality of the spin-helix can be reversed [12]. A recent time-dependent neutron diffraction study showed that the reversal time for this process is of the order of milliseconds [13].

The influence of magnetic order on the lattice dynamics and, hence, the spin–lattice interaction can be efficiently studied by Raman spectroscopy. In previous publications, various authors have already presented Raman data of MnWO₄ [14, 15] and slightly doped MnWO₄ [16] in a wide temperature range, and could not detect any anomaly close to the magnetic phase transitions. Recently, however, another study was published [17] which addressed this temperature region in more detail. In fact, the appearance of a low-frequency excitation...
was observed, which could be attributed to a two-magnon process. As an extension of this previous study, the present investigation is focused on high-resolution experiments on lattice modes and their variation within the magnetically ordered phases and under applied electric fields using the high-resolution triple additive mode of the Raman spectrometer.

2. Experimental setup

The investigated crystal was grown from the melt using a top seeded technique, which is described in [18]. The mosaicity was less than 0.03°, as obtained by γ-ray diffraction. The magnetic properties of this sample were characterized by SQUID magnetometry, which yielded phase transition temperatures close to the published values ($T_1 = 7.5\,\text{K}$, $T_2 = 12.7\,\text{K}$ and $T_N = 13.6\,\text{K}$).

The Raman spectra were recorded using a Jobin-Yvon T64000 spectrometer in backscattering geometry and an Innova 90C argon ion laser with an excitation wavelength of 514 nm and a power of 30 mW. A CCD camera was used for detection. The spectrometer has a focal length of 640 mm and a triple monochromator with three gratings with 1800 lines per mm. The width of the entrance slit was set to 20 μm. For the present high-resolution study, the spectrometer was used in additive mode, yielding a spectral resolution of about 0.3 cm⁻¹. The reproducibility of the spectra is usually limited by random thermal fluctuations of the monochromator and detector. In order to improve the accuracy, a reference spectrum from a neon lamp was measured after the acquisition of each Raman spectrum. By comparing the wavenumbers of the Raman bands with the tabulated neon peaks, the spectrometer is continuously calibrated, thus correcting any possible fluctuations. This technique allows the reduction of the experimental error to less than ±0.1 cm⁻¹ for the absolute wavenumber of a band and to ±0.02 cm⁻¹ for changes during a series of measurements.

The sample was cooled by a Sumitomo SHI-RDK-2025D closed-cycle cryostat. The heat input of the laser was measured with a Si diode next to the specimen to make sure that the temperature is accurately controlled even at very low temperatures. The spectra were obtained in a temperature range from 100 to 6 K. Each phonon was measured without and with an electric field of 6 KV cm⁻¹ applied along the $a$-axis. The incident laser beam was parallel to the $a$-axis and polarized spectra of the configurations $a(bb)a$, $a(bc)a$, $a(cc)a$ and $a(cb)a$ were collected.

3. Results and discussion

3.1. General features of Raman spectra at low temperatures

The unit cell of MnWO₄ contains two formula units and consequently 12 atoms, which lead to 36 vibrational degrees of freedom. The 33 optical modes are divided amongst the four irreducible representations of the point group $P2/m$ according to [15, 16]

$$\Gamma = 8\,A_g + 10\,B_g + 7\,A_u + 8\,B_u.$$  

Being invariant to the inversion, the representations $A_g$ and $B_g$ are Raman active while the representations $A_u$ and $B_u$ are IR active. The $A_g$ modes are visible with parallel polarizers ($a(cc)a$, $a(bb)a$) and the $B_g$ modes appear in configurations with crossed polarizers ($a(bc)a$, $a(cb)a$).

Figures 1 and 2 display all $A_g$ and $B_g$ phonons, respectively. Due to the restricted spectral range covered in additive mode, both figures are composed of several single spectra, as indicated by the vertical dashed lines. The individual spectra were obtained using different acquisition times with ratios indicated by the given amplification factors. For example, $A_g0$ has in fact a one order of magnitude higher intensity than $A_g7$. Note, that several $A_g$ modes are also visible in the cross polarized configuration (figure 2) due to polarization leaks.

As an example for the temperature dependence of Raman modes, selected spectra of $A_g7$ observed in $a(bb)a$ polarization are shown in figure 3. With decreasing temperature the peak narrows and shifts towards higher wavenumbers, just as is usually expected. For most other phonons, however, the changes with temperature are quite subtle and need a careful quantitative evaluation. This was achieved by fitting a Voigt profile, a convolution of a Gaussian and a Lorentzian profile, to each vibrational band. The Gaussian represents the experimental resolution with a full width at half maximum of 0.3 cm⁻¹. The individual phonons are characterized by peak position and Lorentzian width, and will be discussed in the following subsections.

1. Note that for simplicity $a = b \times c$ is used, which is slightly different from the monoclinic lattice vector.
2. Note that the slight asymmetry of this extremely narrow line is due to the experimental conditions at the ultimate resolution limit of the spectrometer [19].
3.2. Temperature dependence of phonon wavenumbers

Figures 4 and 5 show the temperature dependence of the wavenumber for each individual phonon. The vertical solid lines indicate the Néel temperature $T_N = 13.5$ K. The error for the values is $\pm 0.02$ cm$^{-1}$, except for Bg0, Bg6, Bg7 and Bg8. These peaks are of such low intensity that larger fitting errors are obtained.

In general, the temperature range can be divided into two parts, the magnetically ordered phases below $T_N$ and the paramagnetic phase above $T_N$. In the latter, most phonons exhibit normal anharmonic behaviour corresponding to a continuous hardening with decreasing temperature until the harmonic limit is reached. The modes Ag6 and Bg6 are characteristic examples of this behaviour. The wavenumber of Ag3, however, is almost temperature independent between 100 K and $T_N$. A considerable softening with decreasing temperature is observed for the phonons Ag0, Ag1 and Bg1. This finding is consistent with data of [15] and [16], which demonstrate that this softening starts as early as about 180 K. Furthermore, the wavenumbers of the three modes Ag7, Bg5 and Bg8 reach a maximum around 50 K and then start to soften with decreasing temperature.

Entering the magnetically ordered phase, only a few modes are not affected by the transition (Ag1, Ag2, Bg2), while others change their temperature behaviour drastically: there are modes which soften on cooling (Ag0 and Ag1) and others which harden (Ag4, Ag6, Ag7, Bg5, Bg6, Bg8, Bg9) below $T_N$.

Figure 1. Additive spectrum of all Ag phonons composed of four single spectra, as indicated by the dashed lines. The intensities of two bands are comparable only within one single spectrum. The multipliers denote the intensity ratios between the single spectra.

Figure 2. Additive spectrum of all Bg phonons composed of four single spectra, as indicated by the dashed lines. The intensities of two bands are comparable only within one single spectrum. The multipliers denote the intensity ratios between the single spectra. Ag modes that become visible due to the finite efficiency of polarization filters are labelled with an asterisk.
Figure 3. Temperature dependence of the $A_g^7$ band measured with $a(bb)a$ polarization.

Figure 4. Temperature dependence of the wavenumber for all $A_g$ phonons between 100 and 6 K. The vertical line indicates the Néel temperature $T_N = 13.5$ K.
In order to examine the influence of the magnetic phase transitions in more detail, the data for Bg6, Bg5 and Ag7 are shown in figure 6 for the temperature interval between 6 K and 20 K. Apparently, the hardening of these modes starts at $T_2$ rather than at $T_N$, and might therefore be attributed to the multiferroic AF2 phase. In contrast, Ag0 exhibits a strong softening (see figure 7). However, there is no distinct temperature for the onset of the softening, because a precursor effect is observed even well above $T_N$. A linear extrapolation of the behaviour for temperatures above and below $T_N$ leads to an intersection at $T = 12.45 \, K \approx T_2$ (black solid lines in figure 7). This supports the statement that the phonon anomalies are associated with the AF2 phase.

Obviously, the wavenumber variations below $T_2$ cannot be explained by intrinsic anharmonicity. The unusual anisotropic thermal expansion of the crystal lattice [20] might provide another explanation. While the lattice contracts in the $a$- and $c$-directions, there is an expansion in the $b$-direction. In principle, this could explain the different behaviours of phonon modes with their individual eigenvectors. However, the relative
volume change is only of the order of $10^{-5}$ [20], while the relative frequency variations are two orders of magnitude higher for some modes. Hence, the thermal expansion is too weak to explain the phonon anomalies.

A coupling of phonons to the spontaneous polarization could also account for the observed effects. The polarization is caused via inverse Dzyaloshinski–Moriya interactions by the displacement of oxygen atoms which bridge the magnetic manganese atoms. This changes the interatomic distances and could thereby affect the lattice dynamics. However, the polarization vanishes at the lock-in transition at $T_1$. This implies that the wavenumber shifts should be reversed at this temperature, which is not observed at all. Furthermore, all Raman spectra were also measured with an electric field applied along the $b$-axis, parallel to the spontaneous polarization vector. If the distortion of the oxygen atoms influences the phonon frequencies, this effect should be enhanced by an electric field. No differences between the spectra with and without electric field were observed though. The displacement is obviously too small to be detected by Raman spectroscopy. This is also reflected by the weak polarization ($60 \, \mu \text{C m}^{-2}$) of MnWO$_4$ [6].

Another explanation is spin–phonon-coupling, since the onset of a long-range magnetic order has been reported to influence phonon frequencies in particular systems [21–26]. Some of these examples also show changes above $T_N$ due to magnetic fluctuations. Slight shifts of the wavenumber of IR-active modes below $T_N$ have also been reported for MnWO$_4$ and doped MnWO$_4$ [16, 27].

Spin–phonon-coupling occurs when the exchange integral $J(r)$ between two magnetic atoms with the distance $r$ significantly changes during the vibration. According to [24] and [25], the frequency shift $\Delta \omega$ is described as

\begin{align*}
\Delta \omega &= \int J(r) \sin(\omega t) \, dr,
\end{align*}

where $J(r)$ is the exchange integral as a function of the distance $r$. The integral is taken over the entire interatomic distance $r$. For small movements, the integral can be approximated as a linear function of $r$.

**Figure 6.** Variation of the wavenumber for phonons $B_{g6}, B_{g5}$ and $A_{g7}$ between 6 K and 20 K. For the sake of comparison, the values at 12.5 K are subtracted and the individual plots are shifted vertically by 0.05 cm$^{-1}$.

**Figure 7.** Wavenumber of the phonon $A_{g0}$ between 6 K and 20 K. The black solid lines correspond to linear fits from 6 K to 11 K and from 22.5 K to 40 K, respectively. The intersection of the two lines is $T = 12.45$ K.
\[ \Delta \omega = \frac{1}{2 \mu_j \omega_j} \sum_i \frac{\partial^2 F(\mathbf{r})}{\partial u_i^2} \langle S_i \cdot S_{i+r} \rangle \]

where \( \omega_j \) is the frequency of phonon \( j \), \( \mu_j \) the effective mass, \( u_i \) the atomic displacement and \( \langle S_i \cdot S_{i+r} \rangle \) the spin–spin correlation function. The sum runs over all neighbouring magnetic atoms with a non-vanishing exchange integral. In an incommensurate magnetic phase with modulation vector \( \mathbf{k} \), the spin–spin correlation function is replaced by [25]

\[ \langle S_i \cdot S_{i+r} \rangle = K(T) \cos(2\pi \cdot k \cdot r) \]

with a temperature dependent prefactor \( K(T) \). The frequency shift can thereby be rewritten as

\[ \Delta \omega = \frac{K(T)}{\mu_j \omega_j} D_{\text{eff}} \]

with the mode-specific effective spin–phonon-coupling parameter \( D_{\text{eff}} \):

\[ D_{\text{eff}} = \sum_r \frac{\partial^2 F(\mathbf{r})}{\partial u_j^2} \cos(2\pi \cdot k \cdot r). \]

By employing the molecular field approximation, \( K(T) \) and \( \Delta \omega \) become proportional to the square of the sublattice magnetization \( M_{\text{subl}} \). For LaMnO\(_3\) and PrMnO\(_3\), this has in fact been observed [24–26], since both \( M_{\text{subl}} \) and \( \Delta \omega \) are found to vary proportionally to \( \sqrt{T_N - T} \). In the case of MnWO\(_4\), data from polarized neutron diffraction reveal the same behaviour for the intensity of the magnetic superlattice reflection (figure 1(d) of [28]), which is proportional to \( M_{\text{subl}}^2 \). Our results for the wavenumber shifts, however, clearly show a linear temperature behaviour, as illustrated in figures 6 and 7. This indicates that the spin–phonon coupling in this compound is a higher-order phenomenon dominated by the fourth power of \( M_{\text{subl}} \). Consequently, the magnetomechanical or magnetoelectric interactions are quite weak, and both the variation of phonon frequencies and the spontaneous polarization within the multiferroic phase are much smaller compared to those of other manganites.

The variation of the magnetization within the AF3 phase is rather small, as observed by Finger et al [28]. Hence, only minor shifts of phonon frequencies are expected, which are even below the resolution limit for most of the modes. As an exceptional case, the \( A_0 \) mode, which exhibits the strongest softening of all modes, already clearly shows a pronounced precursor effect at temperatures above \( T_N \) (see figure 7). This finding indicates that there is also a weak but significant coupling to magnetic fluctuations. Recent dielectric experiments by Niermann et al [29] seem to support this interpretation. The observed tiny dielectric anomalies close to the transition into the ferroelectric phase led these authors to conclude that there is a softening of low-frequency magnetic excitations with small spectral weight.

In conclusion, we believe that all of the phonon anomalies are due to the spin–phonon coupling and depend on the overall magnetization rather than the details of the modulated structures. In particular, no significant peculiarities are observed at the lock-in transition near 7.5 K.

Characteristic data of all Raman modes are collected in table 1. Column 2 contains the wavenumbers at 10 K while in column 3 their variation within the magnetically ordered phases is listed. It is apparent that, independent of their symmetry, low-frequency phonons, which are essentially determined by the vibration of heavy metal ions, predominantly harden on cooling. In contrast, high-frequency modes, that are dominated by oxygen stretching vibrations, soften considerably. This is illustrated in figure 8. Similar effects have also been observed in ferromagnetic systems such as CdCr\(_2\)S\(_4\) [30] as well as in EuMnO\(_3\), where the Raman–active phonon with the highest frequency exhibits the strongest shift in wavenumber [31], just like \( A_0 \) in our case.

### 3.3. Lorentzian linewidth

The temperature dependence of the Lorentzian linewidth of \( A_0 \) is shown in figure 9 as an example for the general trend for all phonon modes. The linewidths show normal anharmonic behaviour which means they reach a constant value at low temperatures and increase on heating. The individual phonons differ in the low-temperature limit and the magnitude of the temperature coefficient. Both values are listed in table 1, columns 4 and 5, respectively. The accuracy for the determination of the linewidth is in general better than 0.1 cm\(^{-1}\), except for the very weak \( B_0 \) mode. Interestingly, the magnetic phase transitions do not affect the lifetime of any phonon. Rather, the linewidths of all phonons are nearly constant below 30 K, and significant broadening occurs only at temperatures clearly above \( T_N \). This confirms the interpretation in section 3.2 that the wavenumber shifts in the magnetically ordered phases cannot simply be caused by intrinsic anharmonicity.

The low-temperature limits of the linewidths of all Raman modes are plotted versus their wavenumber in figure 10. Obviously, the linewidth of low-energy modes is well below 0.5 cm\(^{-1}\). Only for modes with wavenumbers above 200 cm\(^{-1}\) is a significantly reduced lifetime observed. Since high-energy phonons usually
provide more possible decay channels, the linewidth increases with wavenumber. As indicated by the dashed blue lines in figure 10, however, there are two groups of phonons that are characterized by different slopes of the linear increase of the linewidth. Interestingly, this classification is independent of the irreducible representation of the phonons, as indicated by the different symbols used in figure 10.

The temperature dependence of the linewidth as shown in figure 9 can be explained by anharmonic phonon–phonon interaction based on three-phonon processes where a zone centre phonon decays into two modes with different wavenumbers $\nu_1$ and $\nu_2$ and opposite wavevectors. According to [32] the temperature dependence of $\Gamma$ can be represented by

$$
\Gamma = \Gamma^0 \left( 1 + \frac{1}{\exp\left( \frac{hc\nu_1}{k_B T} \right) - 1} + \frac{1}{\exp\left( \frac{hc\nu_2}{k_B T} \right) - 1} \right). \tag{5}
$$

Here $h$ and $k_B$ are the Planck and Boltzmann constants, respectively, $c$ is the speed of light and $T$ is the temperature. $\Gamma^0$ is the asymptotic value of the linewidth at zero temperature. Taking into account that the sum

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Table 1. Wavenumber of all phonons at 10 K (second column), temperature variation of the wavenumber below $T_N$ (third column), Lorentzian linewidth at 10 K (fourth column) and increase of the Lorentzian linewidth from 6 to 100 K (fifth column).

| Phonon | $\nu$ (10 K) [cm$^{-1}$] | $\nu$ (6 K) − $\nu$ (13.5 K) [cm$^{-1}$] | $\Gamma$ (10 K) [cm$^{-1}$] | $\Gamma$ (100 K) − $\Gamma$ (6 K) [cm$^{-1}$] |
|--------|-------------------------|---------------------------------|-----------------|---------------------------------|
| $A_1^0$ | 884.4 ± 0.1 | -0.26 ± 0.02 | 3.97 ± 0.04 | 0.45 |
| $A_1^1$ | 697.9 ± 0.1 | -0.09 ± 0.02 | 6.78 ± 0.07 | 0.56 |
| $A_2$ | 549.2 ± 0.1 | 0.02 ± 0.02 | 4.66 ± 0.04 | 1.58 |
| $A_3$ | 398.6 ± 0.1 | 0.03 ± 0.02 | 1.58 ± 0.04 | 0.41 |
| $A_4$ | 329.3 ± 0.1 | 0.12 ± 0.02 | 2.40 ± 0.12 | 1.15 |
| $A_5$ | 260.9 ± 0.1 | 0.07 ± 0.02 | 0.75 ± 0.09 | 0.54 |
| $A_6$ | 210.0 ± 0.1 | 0.10 ± 0.02 | 0.94 ± 0.02 | 0.84 |
| $A_7$ | 131.9 ± 0.1 | 0.12 ± 0.02 | 0.40 ± 0.10 | 0.28 |
| $B_1^0$ | 775.5 ± 0.1 | -0.08 ± 0.02 | 6.7 ± 0.3 | 0.96 |
| $B_1^1$ | 673.9 ± 0.1 | 0.00 ± 0.02 | 6.42 ± 0.13 | 0.44 |
| $B_2$ | 514.1 ± 0.1 | 0.04 ± 0.02 | 2.05 ± 0.05 | 0.62 |
| $B_3$ | 460.1 ± 0.1 | -0.02 ± 0.02 | 2.72 ± 0.09 | 0.92 |
| $B_4$ | 296.2 ± 0.1 | 0.07 ± 0.02 | 1.03 ± 0.07 | 0.46 |
| $B_5$ | 275.5 ± 0.1 | 0.16 ± 0.02 | 1.06 ± 0.10 | 0.78 |
| $B_6$ | 180.2 ± 0.1 | 0.14 ± 0.02 | 0.40 ± 0.03 | 0.38 |
| $B_7$ | 168.3 ± 0.1 | 0.09 ± 0.02 | 0.42 ± 0.07 | 0.21 |
| $B_8$ | 162.5 ± 0.1 | 0.20 ± 0.02 | 0.47 ± 0.12 | 0.39 |
| $B_9$ | 91.1 ± 0.1 | 0.07 ± 0.02 | 0.37 ± 0.05 | 0.17 |
and $2\tilde{n}$ is equal to the wavenumber of the excited phonon, we were able to fit the temperature dependence of the linewidth, as exemplified in figure 9.

The fit result for the lower wavenumber $\tilde{v}_2$, which characterizes the dominant term in equation (5) at sufficiently high temperatures, is around $\tilde{v}_2 \approx 100$ cm$^{-1}$ for most of the phonons. In view of the observed optical modes (see table 1 and [27]), it is expected that the first maximum of the phonon density of states appears just in this energy range. Consequently, considerable decay can only be expected for phonons with wavenumbers above $2\tilde{n}$. This is in agreement with the observation that a distinct threshold for broadening exists at about 200 cm$^{-1}$, as shown in figure 10. Only for the high-frequency stretching modes above approximately 670 cm$^{-1}$ ($A_{g0}, A_{g1}$, and $B_{g1}$) are the decay channels dominated by higher-energy modes corresponding to larger values of $\tilde{v}_2$. This finding is consistent with the observation that the phonon density of states exhibits a significant gap between 550 cm$^{-1}$ and 670 cm$^{-1}$.

4. Conclusion

Using Raman spectroscopy with extremely high precision, we were able to observe a variety of phonon anomalies that are associated with the magnetic order in MnWO$_4$ below $T_N = 13.5$ K. The employment of the triple additive mode accompanied by careful and continuous calibration procedures allowed the determination of phonon wavenumber with an accuracy of 0.02 cm$^{-1}$. It was found that low-energy modes become harder on...
cooling below $T_c$, while the metal–oxygen stretching vibrations soften. The most pronounced effect of more than 0.3 cm$^{-1}$ was observed for the highest-frequency $A_0$ mode. Detailed investigations of the temperature dependence show an almost linear behaviour for all modes and provide evidence for higher-order spin–phonon coupling. Other possible explanations such as the effect of anharmonicity or anomalous thermal expansion are discussed and can be ruled out. Moreover, no effect of the electric polarization within the multiferroic phase could be identified.

From the results of phonon linewidths we inferred that no significant damping is present below about 30 K, and consequently magnetic order has no influence on the phonon lifetime. Only for temperatures well above any magnetic transition is the usual anharmonic behaviour found, as described by phonon–phonon interactions. The higher the phonon frequency the larger the number of possible decay channels and hence the phonon linewidth.

Despite the weak magneto-electric coupling in MnWO$_4$, we were able to identify significant anomalies in the phonon spectra caused by the magnetic ordering process and provide new information about the underlying microscopic mechanism.

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