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Abstract

We propose and experimentally demonstrate a means of broadband phonon-magnon interconversion that relies on combining magnetoelastic coupling with translational symmetry breaking in the important experimental material yttrium iron garnet (YIG). As well as being of interest for its basic physics, this quasiparticle coupling mechanism adds to the range of effects that potentially find useful application in hybrid solid-state quantum computing devices as well as low-power wave-based classical computing architectures.

The magnon is a relative newcomer to the catwalk of hybrid solid-state quantum science but it has already begun to turn heads. The rich physics and ready tunability of microwave magnonic systems combined with their demonstrable compatibility with the existing tools of experimental solid-state quantum engineering suggest significant and wide-reaching opportunities, not only in device design, but also in fundamental research [1–14]. Moreover, in the context of classical computing, there has been steadily growing interest in the use of magnonic systems as a platform for wave-based information technologies that overcome the ever more pressing ‘heat death’ issues associated with conventional computing hardware [15–28].

Indeed, phase-modulated spin-waves have significant appeal as data carriers in both classical and quantum computing devices: they offer Joule-heat-free spin information transfer and their short wavelengths relative to electromagnetic waves of the same frequency (microwave-frequency spin waves have wavelengths in the millimetre to nanometre range) are highly conducive to progressive device miniaturization [24–28]. Moreover, the ability to interconvert between spin-wave or magnon signals and those in other physical domains—notably microwave photonics, spin currents, heat currents, and optics—is widely recognised as a further important dividend [27–30].

However, until now, a notable gap has existed in the catalogue of magnonic conversion effects. Though the coupling between the magnon and phonon systems of magnetic materials was, in fact, the inspiration for the original theoretical framework upon which all of spin-wave and magnon physics came to be based [31], the interconversion between magnon and phonon *signals*—as opposed to incoherent excitations—had yet to be practically demonstrated [32, 33].

In this paper, we propose and demonstrate the first experimental proof of principle of a novel phonon-based approach to magnon signal generation. The effect is predicated on a new quasiparticle coupling mechanism with two essential ingredients: magnetoelastic coupling of sufficient strength and appropriate symmetry in the magnonic host material; and energy-momentum matching between the phonons and the magnons. The latter is generally difficult to realise since the phonon and magnon dispersion relations overlap and hybridise only over a very narrow range of wavenumbers that would be impractical for broadband signal transfer. As explained below, we circumvent this by borrowing the artifice of translational symmetry breaking that is used to solve similar issues in other areas of wave physics.

The key requirement of magnetoelastic coupling is present to exactly the right degree in the magnetic material we choose for our experiments, the popular electrically insulating ferrimagnet yttrium iron garnet or YIG ($\text{Y}_3\text{Fe}_5\text{O}_{12}$). The magnetoelastic coupling in YIG is of sufficient magnitude to provide adequate phonon–magnon mode coupling, yet insufficiently strong to compromise the spin-wave decay length. These favourable properties arise on account of the fact that YIG’s magnetic moment comes from S-state Fe^{3+} ions [34] and that, as is typical of magnetic garnets, the chemical and magnetic unit cells are inequivalent [35].

To address the energy/wavenumber (k) matching issue we take our cue from the operating principle of a Yagi television aerial which phase-matches guided and free-space electromagnetic waves with very different dispersions. The Yagi does this through the expedient of breaking translational symmetry by establishing standing-wave patterns with a geometry calculated to generate spatial frequencies in the local electromagnetic disturbance that match the wavenumbers of the quasiparticles that it is required to produce. We mimic this approach by using a piezoelectric transducer to send a narrow pulsed microwave acoustic beam which impinges locally on a magnonic waveguide fashioned from a YIG film that is subject to an externally applied magnetic field. This process generates a small magnetoelastically active volume defined by the beam profile and the waveguide thickness. The local symmetry break due to the presence of the waveguide interfaces and the resultant acoustic standing wave structure allows the low k magnons to couple effectively to phonons of much higher k values, i.e. well removed from the usual hybridisation regime. Since our aim is a very broadband phonon–magnon conversion process, we generate a spatial frequency spectrum that is as white as possible by engineering the waveguide and the beam to confine the region of magnetoelastic interaction as tightly as possible in all three dimensions.

Due to the magnetoelastic coupling, the acoustic displacement of the YIG film locally generates a small signal oscillating internal magnetic field contribution $\mathbf{h}_{\text{me}}$, some of whose spatial frequency components are capable of coupling to and exciting MSW modes. In principle, $\mathbf{h}_{\text{me}}$ couples to an infinite family of propagating modes; however, excitation of the higher order modes is inefficient owing to their low propagation speeds, and hence only the lowest order mode is excited. For a given magnetoelastic free energy density F_{me} , the oscillating magnetic field is

$$\mathbf{h}_{\text{me}} = -\frac{1}{M_s} \frac{\partial F_{\text{me}}}{\partial \boldsymbol{\alpha}}. \quad (1)$$

Since the magnetic structure of YIG is effectively cubic, the phenomenological magnetoelastic free energy density F_{me} is given by [36]:

$$F_{\text{me}} = B_1 \sum_i \alpha_i^2 \epsilon_{ii} + B_2 \sum_{i \neq j} \alpha_i \alpha_j \epsilon_{ij}, \quad (2)$$

where B_1 and B_2 are the first order magnetoelastic coefficients, $\alpha_i = \frac{M_i}{M_s}$ is the directional cosine with respect to the i th cubic axis, M_i is the magnetization component in the i -direction, M_s is the saturation magnetization and ϵ_{ij} is the strain tensor. For a $\langle 111 \rangle$ oriented YIG crystal with longitudinal acoustic excitation, the effective magnetic field $\mathbf{h}_{\text{me}}$ can be calculated using equations (1) and (2) such that [37, 38]:

$$\mathbf{h}_{\text{me}} = -\frac{1}{M_s} B_2 \sin(2\theta) \epsilon_{33} \hat{\mathbf{y}} \quad (3)$$

where ϵ_{33} represents the strain caused by the longitudinal acoustic wave, $\hat{\mathbf{y}}$ is the unit vector pointing perpendicular to the net magnetization direction and the $\langle 11\bar{2} \rangle$ axis of the YIG film, θ is the angle made by the total internal field with the surface normal (the internal field direction is determined by the combination of the externally applied bias field and the demagnetising field: see figure 1 and supporting text).

Since the electrical impedance of the acoustic transducer used is intrinsically high, the result is a magnon source that creates minimal Ohmic heating and that is not only broadband but also capable of highly spatially localised magnonic excitation.

Our acoustic magnon launcher is illustrated schematically in figure 1. A waveguide that supports spin-wave transmission is fashioned from a $7.8 \mu\text{m}$ thick monocrystalline YIG film grown by liquid phase epitaxy on a 0.45 mm thick gallium gadolinium garnet (GGG) substrate. The surface normal of the waveguide is in the direction of the $\langle 111 \rangle$ crystal axis of the YIG; the $\langle 1\bar{1}0 \rangle$ axis is parallel to the waveguide axis, and therefore to the direction of measured MSW propagation.

The magnon waveguide is placed in an external bias field $\mathbf{B}_{\text{ext}}$ applied at an angle θ_{ext} to the surface normal. The field is oriented in the plane containing the surface normal and the propagation direction, and in this configuration, volume MSW modes can be excited (i.e. those whose magnetic amplitude extends throughout the volume of the film) [25, 26]. At $\theta_{\text{ext}} = 0^\circ$, only forward volume magnetostatic spin-wave (FVMSW) modes can propagate (group velocity parallel to phase velocity). In contrast, when $\theta_{\text{ext}} = 90^\circ$, the waveguide can only support backward volume magnetostatic spin-wave (BVMSW) modes (group velocity antiparallel to phase velocity). When the magnetic field is pointing at an arbitrary angle θ_{ext} to the normal of the plane, FVMSW and BVMSW modes are excited simultaneously.

Grown on the back (substrate) side of the waveguide is an array of zinc oxide (ZnO) ultrasonic transducers whose circular conductive top contacts are spaced by 0.5 mm along the waveguide axis. The transducer thickness is of order 100 nm , (this corresponding to the transducer functioning approximately as a quarter wave dipole in the microwave frequency band of interest). The transducers are capable of injecting longitudinally polarised bulk acoustic wave (BAW) beams at GHz frequencies perpendicular to the YIG surface. Their

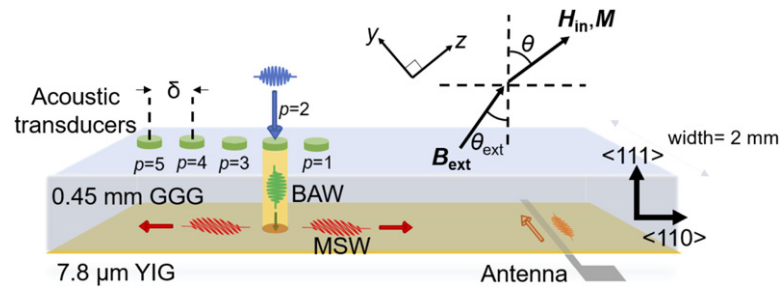


Figure 1. Schematic of the acoustic spin-wave excitation experiment. An array of ZnO transducers (diameter $300\ \mu\text{m}$, spacing $\delta = 0.5\ \text{mm}$) is deposited on the reverse of the $0.45\ \text{mm}$ thick GGG substrate to which a YIG waveguide is attached. The YIG is $7.8\ \mu\text{m}$ thick and has its $\langle 111 \rangle$ axis in the direction of the waveguide normal. An inductive microwave output antenna for the detection of MSWs is placed a known distance from the transducer array [$d_p = (0.86\ \text{mm} + (p-1) \times 0.5\ \text{mm})$]. An external magnetic bias field B_{ext} is applied at an angle θ_{ext} with respect to the normal of the waveguide. The field is oriented in the plane containing the surface normal and the propagation direction. The internal field H_{in} (which is parallel to the magnetisation vector M) points at an angle $\theta \geq \theta_{\text{ext}}$, due to the effect of the demagnetising field.

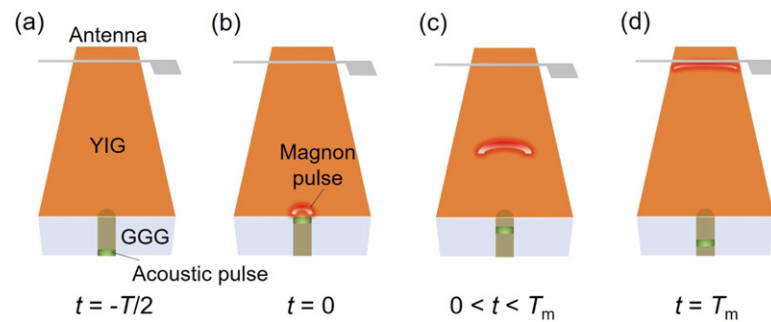
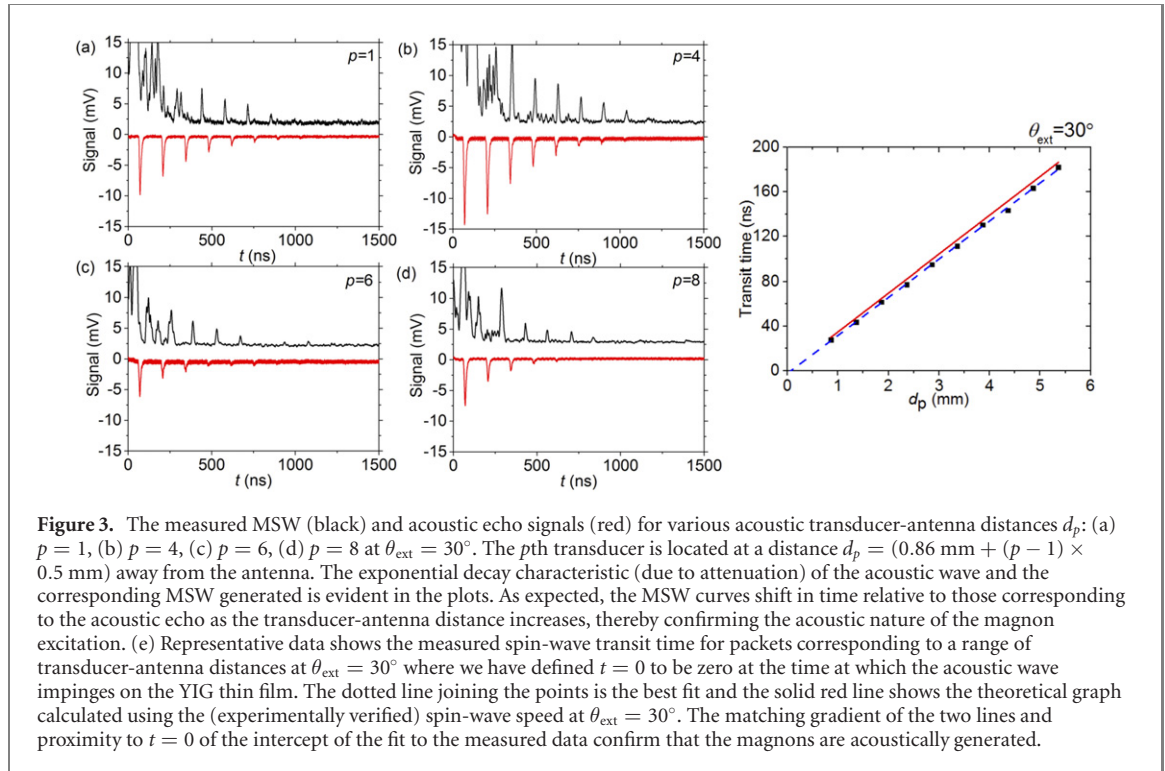


Figure 2. The sequence of the experiment from the launching of an acoustic pulse to the reception of the MSW packet at the microwave output antenna. T = acoustic round trip time, T_m = magnon transit time. (a) At $t = -T/2$, the acoustic wave is launched. (b) At $t = 0$, the acoustic wave impinges onto the YIG film. (c) The acoustics generate a laterally propagating MSW. (d) At $t = T_m$ the magnon pulse arrives at the microwave output antenna.

diameter is chosen to be $300\ \mu\text{m}$ as a compromise between generating a narrow pencil beam and minimising beam walkoff due to diffraction effects and surface imperfections. Electrical contacts to the transducers are made using a microwave-frequency probe with a tip diameter of $125\ \mu\text{m}$. When the acoustic transducer is excited, the resulting pulse of acoustic energy propagates across the thick GGG substrate and impinges on the interface with the comparatively thin YIG film. The acoustic impedance of YIG is very similar to that of the GGG [39, 40], so only a negligibly small fraction (around 1%) is reflected from the YIG/GGG interface. The acoustic wave then propagates through the YIG before it encounters the interface between the YIG and the surrounding air. Since the acoustic impedance of air is very low in comparison with that of the YIG, all of the acoustic energy is reflected and the acoustic amplitude is maximum at the surface. As the reflected wave propagates back through the thickness of the system it interferes with the incoming signal that created it, forming a standing wave.

To demonstrate that magnons can be excited by the spatially localised phonons generated by the acoustic source, we use a single inductive antenna to monitor the spin-wave propagating through the waveguide when it is subject to an acoustic pulse. The antenna is positioned at one end of the waveguide, separated from the closest end of the acoustic transducer array by a distance of $0.86\ \text{mm}$. The sequence of events that occur after the acoustic pulse is launched is illustrated schematically in figure 2. Firstly, at $t = -T/2$, where T is the round trip time of the acoustic wave, a BAW pulse is launched by applying a $15\ \text{ns}$ electrical pulse with a $3\ \text{GHz}$ carrier frequency to the top electrode of a ZnO transducer. After time $T/2$ ($t = 0$), the acoustic pulse hits the YIG surface, thereby creating a lateral propagating spin-wave. This wave propagates along the waveguide, and at $t = T_m$, arrives at the inductive antenna. When the acoustic wave returns into the transducer at $t = T/2$, an acoustic echo signal, reduced in amplitude relative to the original due to the round-trip attenuation, is detected via the piezoelectric effect. This cycle repeats for as long as the acoustic wave has sufficient amplitude to be detectable. In our experiments, both the magnon signal detected at the electrical antenna and the acoustic



echo signal detected by the piezoelectric transducer are amplified, rectified, and simultaneously recorded using a 10 GHz sampling oscilloscope.

We first measure the magnetic response of the waveguide to an acoustic pulse generated by the transducer closest to the inductive output antenna: a distance of 0.86 mm. The external bias field angle θ_{ext} is set to 30° and the field strength is tuned such that a 3 GHz spin-wave signal sits within the passband of the YIG waveguide.

Figure 3(a) shows experimental data from the system's output antenna corresponding to such a configuration with $B_{\text{ext}} = 90 \text{ mT}$. The acoustic echo signals are plotted in the same graph. The noise near the point when the electromagnetic signal was injected initially is caused by the excitation of spin-waves by the stray electromagnetic field of the probe and direct electromagnetic coupling with the antenna. Some considerable time after the initial excitation, we see that for every acoustic echo, there is a corresponding magnon signal. As expected, the acoustic wave decays exponentially due to the acoustic attenuation of the medium in which it is travelling and the corresponding magnon signal decays proportionally. We measure the magnon transit time to be 27.8 ns by finding the difference in time between the mid-point of two consecutive acoustic echoes (i.e. the time at which the phonons hit the YIG) and the arrival time of the magnon signal. This figure is in line with the predictions made both by earlier experimental measurements and by direct calculation of the magnon velocity [41–43].

To confirm the acoustic origin of the magnon signal we repeat the experiment using acoustic transducers along the axis of the waveguide separated from the inductive output antenna by progressively larger distances. The p th transducer is located at a distance $d_p = (0.86 \text{ mm} + (p - 1) \times 0.5 \text{ mm})$ from the output antenna. Conducting experiments using transducers for all values of p we observe that the corresponding magnon peak arrives after a transit time that scales linearly with d_p with a gradient that agrees closely with the known value of magnon velocity (figure 3). Furthermore, the y -intercept of the plot is very close to zero (the discrepancy is only -2.5 ns which is within experimental measurement error), confirming beyond reasonable doubt that the spin-waves are indeed generated as the acoustic waves hit the YIG.

As predicted in equation (3), the efficiency of the acoustic/MSW coupling mechanism is strongly dependent on the orientation of the acoustic and waveguide axes with respect to the applied external magnetic field. Plotted in figure 4 is the experimentally obtained phonon–magnon conversion efficiency as a function of the angle of the externally applied magnetic field θ_{ext} (black squares). Overlaid is a solid line showing the form of the data predicted by our theoretical model. Naively, one might expect the efficiency curve to follow the $\sin(2\theta)$ dependence of equation (3) but the reality is more complex because the data is the product of no less than four separate effects: the $\sin(2\theta)$ dependence of the conversion process; the angular dependence of the coupling efficiency of the receiving antenna; the angular dependence of the spin-wave attenuation which determines the loss in signal between magnon creation and detection; and the relation between the angles of the internal

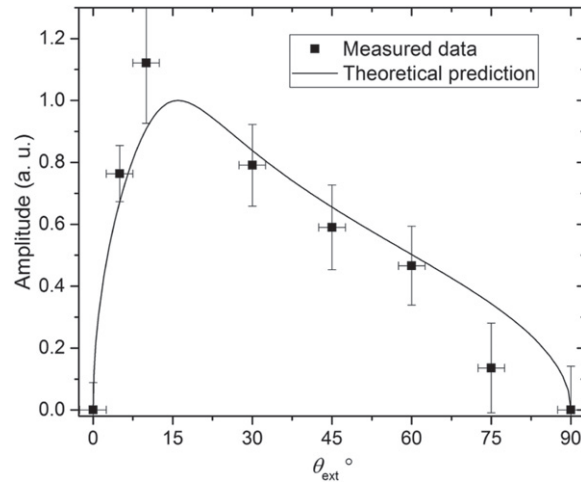


Figure 4. The experimentally obtained phonon-magnon conversion efficiency as a function of the angle, θ_{ext} of the externally applied magnetic field (black squares). The measured amplitude at a field angle of 45° corresponds to a phonon–magnon conversion efficiency of approximately -40 dB. Overlaid is a solid line showing the form of the data predicted by our theoretical model.

magnetic field and the externally applied field. Within the error-bars that arise due to field inhomogeneity and carrier interference, it can be seen that there is good agreement between theoretical model and experimental data. We note that we observe little or no magnon excitation when the external field is pointing normal to or in the plane of the waveguide, whereas a considerable magnon signal amplitude is observed when the field is at an intermediate angle. This result is in accord with the fact that no oscillating magnetic field h_{me} is produced when $\theta_{\text{ext}} = \theta = 0^{\circ}$ or 90° (equation (3)). Under typical operating conditions, the transducers deliver 15 ns pulses with an acoustic power of 1 mW which corresponds to an acoustic strain amplitude of order 3×10^{-6} . These phonon pulses generate magnon pulses with a power of approximately -40 dB m, thereby giving a measured phonon-magnon conversion efficiency for YIG of -40 dB with field angle $\theta_{\text{ext}} = 45^{\circ}$.

The temporal response of the phonon to magnon conversion mechanism has also been simulated and measured as a function of B_{ext} at various θ_{ext} . Our simulator takes an initial time-dependent Gaussian signal $f(t)$ representing the pulse excitation and decomposes it into its Fourier components $g(\omega)$ using FFT. Each component accumulates a phase of $e^{i\omega d/v_p}$ as it propagates over a distance d . Here, v_p denotes the phase velocity which is governed by the dispersion relation of the spin wave [41–43]. The k -dependent coupling factors of the top dot $A_{\text{topdot}}(k)$ and the antenna $A_{\text{antenna}}(k)$ which are defined by their corresponding geometries are also taken into account. These factors are then transformed into functions of ω via the dispersion relation. The output time-dependent signal is eventually computed by inverse Fourier transforming the function $G(\omega) = g(\omega) \cdot A_{\text{topdot}}(\omega) \cdot A_{\text{antenna}}(\omega) \cdot e^{i\omega d/v_p}$. Shown in figure 5 is the temporal evolution of the measured and simulated voltages at the spin-wave detection antenna for a certain range of B_{ext} at $\theta_{\text{ext}} = 10^{\circ}$ and $\theta_{\text{ext}} = 45^{\circ}$. For $\theta_{\text{ext}} = 10^{\circ}$ (figure 5(a) and (b)), two distinct regions of excitation are observed at low and high magnetic fields corresponding to the FVMSW and BVMSW modes respectively. When $\theta_{\text{ext}} = 10^{\circ}$, the internal field is pointing at an angle θ less than 45° with respect to the surface normal, in this regime, the FVMSW is dominant and the BVMSW have comparatively smaller velocity and amplitude. Conversely, when $\theta_{\text{ext}} = 45^{\circ}$, $\theta > 45^{\circ}$, under this condition, the BVMSW dominate so significantly that the FVMSW are too weak to be detected (figure 5(c) and (d)). There is excellent quantitative agreement between our simulated and experimental results. The periodic horizontal feature in the experimental data is an artefact caused by interference of a small leaked portion of the microwave carrier signal with the output signal.

In summary, we have demonstrated for the first time, the use of GHz-frequency piezoelectrically generated phonons for the direct excitation of magnon signals. As well as being of significant interest for its basic physics, this mechanism opens doors to a range of technological opportunities. Though our experiments have been conducted at room temperature in the classical regime, the technique is directly transferable to low-temperature quantum experiments. As such, it represents a bridge between the world of microwave-frequency quantum acoustics [44–49] and microwave-frequency quantum magnonics [1–14], potentially allowing the advantages of these two versatile but very different signalling environments to be combined. The mechanism has a range of features that make it both technologically promising and practically attractive in the context of experiments and device development in the quantum and classical regimes. Firstly, the high impedance piezoelectric acoustic transducers at the heart of the technique have very low power dissipation, particularly when compared with the inductive antennae traditionally employed in the context of room-temperature device

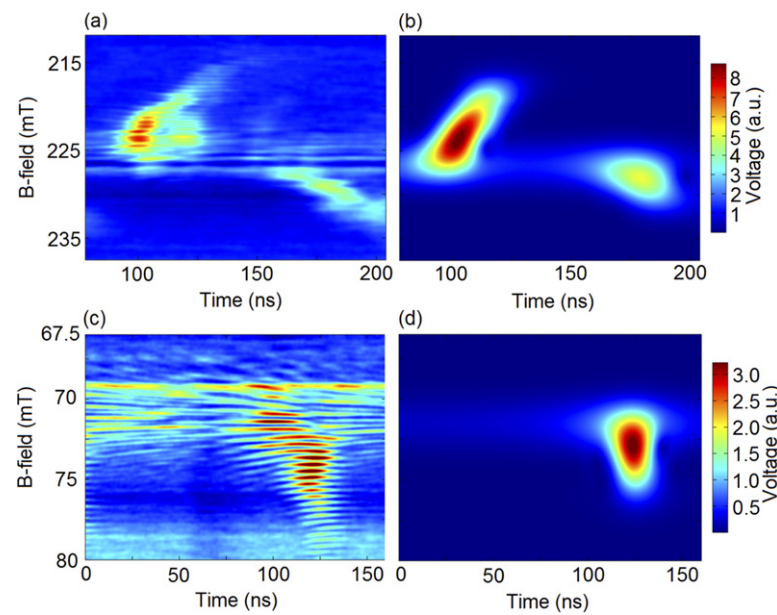


Figure 5. (a) The voltage measured at the inductive antenna as a function of applied magnetic field and time at $\theta_{\text{ext}} = 10^\circ$. FVMSW and BVMSW modes are excited at lower and higher magnetic bias fields respectively. (b) Simulated results for $\theta_{\text{ext}} = 10^\circ$. (c) The voltage measured at the antenna as a function of applied magnetic field and time at $\theta_{\text{ext}} = 45^\circ$. Only BVMSW modes are excited in this case. (d) Simulated results for $\theta_{\text{ext}} = 45^\circ$. The horizontal striation observed in the experimental results is due to the interference caused by a small amount of carrier breakthrough.

design. Secondly, the piezoelectric transducers lend themselves naturally to miniaturisation: lateral patterning and periodic etching can be combined to engineer even shorter phonon wavelength, higher efficiency devices including those that work in conjunction with exchange spin-waves having wavelengths $\ll 1 \mu\text{m}$. Finally, unlike inductive microwave antennas, the acoustic magnon source is point-like, paving the way for entirely new forms of waveform engineering inside magnon circuits. In the context of YIG-based quantum magnonics the work suggests a particularly enticing range of opportunities. As discussed earlier, YIG has recently come under the spotlight in the context of solid-state quantum device design. The material is important not only on account of its very low magnon damping, but also because it potentially offers mechanisms for the interconversion between signals of different character and frequency range—notably, as discussed in reference [29] (and accompanying references) microwave to optical conversion. Our new interconversion process expands these opportunities yet further. Most especially, it suggests the possibility of producing novel quantum acoustic devices for hybrid quantum electrodynamics that exploit the transduction of longitudinal acoustic waves in combination with magnon-photon interconversion, and magnonic devices in which acoustic transduction is used to access magnetic modes that cannot be addressed easily (or at all) electromagnetically. In particular, the acoustic excitation method makes it possible to achieve point excitation of travelling magnon modes in bulk samples: a process that is very difficult to achieve using electromagnetic antennae and most attractive from an experimental point of view.

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Data availability statement

The data that support the findings of this study are available upon reasonable request from the authors.

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