

Controlling acoustic waves using magneto-elastic Fano resonances

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We propose and analyze theoretically a class of energy-efficient magneto-elastic devices for analogue signal processing. The signals are carried by transverse acoustic waves while the bias magnetic field controls their scattering from a magneto-elastic slab. By tuning the bias field, one can alter the resonant frequency at which the propagating acoustic waves hybridize with the magnetic modes, and thereby control transmission and reflection coefficients of the acoustic waves. The scattering coefficients exhibit Breit-Wigner/Fano resonant behaviour akin to inelastic scattering in atomic and nuclear physics. Employing oblique incidence geometry, one can effectively enhance the strength of magneto-elastic coupling, and thus countermand the magnetic losses due to the Gilbert damping. We apply our theory to discuss potential benefits and issues in realistic systems and suggest routes to enhance performance of the proposed devices.

Optical and, more generally, wave-based computing paradigms gain momentum on a promise to replace and complement the traditional semiconductor-based technology.¹ The energy savings inherent to non-volatile memory devices has spurred the rapid growth of research in magnonics,^{2,3} in which spin waves⁴ are exploited as a signal or data carrier. Yet, the progress is hampered by the magnetic loss (damping).^{5,6} Indeed, the propagation distance of spin waves is rather short in ferromagnetic metals while low-damping magnetic insulators are more difficult to structure into nanoscale devices. In contrast, the propagation distance of acoustic waves is typically much longer than that of spin waves at the same frequencies.⁷ Hence, their use as the signal or data carrier could reduce the propagation loss to a tolerable level. Notably, one could control the acoustic waves using a magnetic field by coupling them to spin waves within magnetostrictive materials.^{8–10} To minimize the magnetic loss, the size of such magneto-acoustic functional elements should be kept minimal. This implies coupling propagating acoustic waves to confined spin wave modes of finite-sized magnetic elements. As we show below this design idea opens a route towards hybrid devices combining functional benefits of magnonics^{2,3} with the energy efficiency of phononics.^{7,11,12}

The phenomena resulting from interaction between coherent spin and acoustic waves have already been addressed in the research literature: the spin wave excitation of propagating acoustic waves^{7,13–15} and vice versa,^{8,16–18} acoustic parametric pumping of spin waves,^{19–21} magnon-phonon coupling in cavities,^{22–24} and mode locking,²⁵ magnonic-phononic crystals,^{26,27} Bragg scattering of spin waves from a surface acoustic wave induced grating,^{28–30} topological properties of magneto-elastic excitations,^{15,31} acoustically driven spin pumping and spin Seebeck effect,^{32,33} and optical excitation and detection of magneto-acoustic waves.^{34–40} However, studies of the interaction between propagating acoustic waves and spin wave modes of finite-sized magnetic elements, which are the most promising for applications, have been relatively scarce to date.^{10,34,36,39}

Here, we explore theoretically the class of magneto-

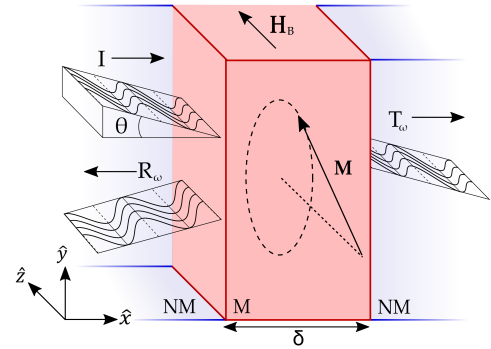


FIG. 1. The prototypical magneto-elastic resonator is a thin magnetic slab (M) of width δ , biased by an external field \mathbf{H}_B , and embedded into a non-magnetic (NM) matrix. The acoustic wave with amplitude I incident at angle θ induces precession of the magnetisation vector \mathbf{M} via the magneto-elastic coupling. As a result, the wave is partly transmitted and reflected, with respective amplitudes T_ω and R_ω .

acoustic devices in which the signal is carried by acoustic waves while the magnetic field controls its propagation via the magnetoelastic interaction in thin isolated magnetic inclusions as shown in Fig. 1. By changing the applied magnetic field, one can alter the frequency at which the incident acoustic waves hybridize with the magnetic modes of the inclusions. Thereby, one can control the acoustic waves by the resonant behaviour of Breit-Wigner and Fano resonances in the magnetic inclusion.⁴¹ We find that the strength of the resonances is suppressed by the ubiquitous magnetic damping in realistic materials, but this can be mitigated by employing oblique incidence geometry. To compare magneto-acoustic materials for such devices, we introduce a figure of merit. The magneto-elastic Fano resonance is identified as most promising in terms of frequency and field tuneability. To enhance resonant behaviour, we explore the oblique incidence as a means by which to enhance the figure of merit.

We consider the simplest geometry in which magneto-elastic coupling can affect sound propagation. A ferromag-

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netic slab ("magnetic inclusion") of thickness δ , of the order of 10nm, is embedded within a non-magnetic medium (Fig.1). The slab is infinite in the $Y - Z$ plane, has saturation magnetization M_s , and is biased by the applied field $\mathbf{H}_B = H_B \hat{z}$. Due to the magneto-elastic coupling, this equilibrium configuration is perturbed by shear stresses in the xz - and yz planes associated with the incident acoustic wave.

To derive the equations of motion, we represent the magnetic energy density F of the magnetic material as a sum of the magneto-elastic F_{ME} and purely magnetic F_M contributions.⁴² Taking into account the Zeeman and demagnetizing energies, we write $F_M = -\mu_0 \mathbf{H}_B \mathbf{M} + \frac{\mu_0}{2} (N_x M_x^2 + N_y M_y^2)$, where $N_{x(y)}$ are the demagnetising coefficients, $N_x + N_y = 1$, \mathbf{M} is the magnetization and μ_0 is the magnetic permeability. In a crystal of cubic symmetry, the magnetoelastic contribution takes the form⁴³

$$F_{ME} = \frac{B}{M_s^2} \sum_{i \neq j} M_i M_j u_{ij} + \frac{B'}{M_s^2} \sum_i M_i^2 u_{ii}, \quad i, j = x, y, z, \quad (1)$$

where B' and B are the linear isotropic and anisotropic magneto-elastic coupling constants, respectively.⁴⁴ The strain tensor is $u_{jk} = \frac{1}{2} (\partial_j U_k + \partial_k U_j)$, where U_j are the displacement vector components. To maximize the effect of the coupling B , we consider a transverse acoustic plane wave incident on the slab from the left and polarized along the bias field, so that $U_x = U_y = 0$, $U_z = U(x, y, t)$. The non-vanishing components of the strain tensor are $u_{xz} = \frac{1}{2} \partial_x U$ and $u_{yz} = \frac{1}{2} \partial_y U$, and F_{ME} is linear in both \mathbf{M} and U :

$$F_{ME} = \frac{B}{M_s} (M_x u_{xz} + M_y u_{yz}). \quad (2)$$

The magnetization dynamics in the slab is due to the effective magnetic field, $\mu_0 \mathbf{H}_{\text{eff}} = -\delta F / \delta \mathbf{M}$. We define \mathbf{m} as the small perturbation of the magnetic order, i.e. $|\mathbf{m}| \ll M_s$. Linearizing the Landau-Lifshitz-Gilbert equation,⁴ we write

$$-\frac{\partial m_x}{\partial t} = \gamma \mu_0 (H_B + N_y M_s) m_y + \gamma B \frac{\partial U}{\partial y} + \alpha \frac{\partial m_y}{\partial t}, \quad (3)$$

$$\frac{\partial m_y}{\partial t} = \gamma \mu_0 (H_B + N_x M_s) m_x + \gamma B \frac{\partial U}{\partial x} + \alpha \frac{\partial m_x}{\partial t}, \quad (4)$$

where γ is the gyromagnetic ratio and α is the Gilbert damping constant. To describe the acoustic wave, we include the magneto-elastic contribution to the stress, $\sigma_{jk}^{(ME)} = \delta F_{ME} / \delta u_{jk}$, into the momentum balance equation:

$$\rho \frac{\partial^2 U}{\partial t^2} = \frac{\partial}{\partial x} \left(C \frac{\partial U}{\partial x} + \frac{B}{M_s} m_x \right) + \frac{\partial}{\partial y} \left(C \frac{\partial U}{\partial y} + \frac{B}{M_s} m_y \right), \quad (5)$$

where $C = c_{44}$ is the shear modulus and ρ is the mass density. The non-magnetic medium is described by Eq.(5) with $B = 0$.

Since the values of C , B , and $N_{x,y}$ are constant within each individual material, we shall seek solutions of the equations in the form of plane waves $U, m_{x(y)} \propto \exp[i(k_{\omega,x} x + k_{\omega,y} y - \omega t)]$. From herein, we consider all variables in the Fourier domain. For the magnetization precession in the magnetic layer driven

by the acoustic wave, we thus obtain

$$m_x = \frac{\gamma B (\omega k_{\omega,y} + i \tilde{\omega}_y k_{\omega,x})}{\omega^2 - \tilde{\omega}_x \tilde{\omega}_y} U, \quad (6)$$

$$m_y = \frac{i \gamma B (\tilde{\omega}_x k_{\omega,y} + i \omega k_{\omega,x})}{\omega^2 - \tilde{\omega}_x \tilde{\omega}_y} U, \quad (7)$$

where we have denoted $\omega_{x(y)} = \gamma \mu_0 (H_B + N_{x(y)} M_s)$ and $\tilde{\omega}_{x(y)} = \omega_{x(y)} - i \omega \alpha$. The complex-valued wave number k_{ω} is given by the dispersion relation

$$k_{\omega}^2 = \frac{\omega^2 \rho (\omega^2 - \tilde{\omega}_x \tilde{\omega}_y)}{C \left[\omega^2 - \tilde{\omega}_x \tilde{\omega}_y + \frac{\gamma B^2}{M_s C} (\tilde{\omega}_y \cos^2 \theta + \tilde{\omega}_x \sin^2 \theta) \right]}. \quad (8)$$

Eq. (8) describes the hybridization between acoustic waves and magnetic precession at frequencies close to ferromagnetic resonance (FMR) at frequency ω_{FMR} , with linewidth Γ_{FMR} . The frequency at which the precession amplitudes (Eqs. (6) and (7)) diverge is given by the condition $(\omega_{\text{FMR}} + i \Gamma_{\text{FMR}}/2)^2 = \tilde{\omega}_x \tilde{\omega}_y$. In the limit of small α , this yields $\omega_{\text{FMR}} = \omega_x \omega_y$ and $\Gamma_{\text{FMR}} = \alpha (\omega_x + \omega_y)$. Away from the resonance, Eq. (8) gives the linear dispersion of acoustic waves. In the non-magnetic medium ($B = 0$), one finds $k_0^2 = \omega^2 \rho_0 / C_0$. Here and below, the subscript '0' is used to mark quantities pertaining to the non-magnetic matrix.

To calculate the reflection and transmission coefficients, R_{ω} and T_{ω} , for a magnetic inclusion, we introduce the mechanical impedance as $Z = i \sigma_{xz} / \omega U_{\omega}$. Solution of the wave matching problem can then be expressed via the ratio of load (Z_{ME}) and source (Z_0) impedances. For the magnetic slab, we find

$$Z_{\omega, \text{ME}} = \frac{C k_{\omega} \cos \theta}{\omega} \left(1 + \frac{\gamma B^2 \tilde{\omega}_y}{M_s C (\omega^2 - \tilde{\omega}_x \tilde{\omega}_y)} \right). \quad (9)$$

For the non-magnetic material, Eq. (9) recovers the usual acoustic impedance⁴⁵ $Z_0 = \cos \theta \sqrt{\rho_0 C_0}$. Due to magnon-phonon hybridization, $Z_{\omega, \text{ME}}$ diverges at ω_{FMR} and vanishes at a nearby frequency ω_{ME} . For $\alpha = 0$, the latter is given by

$$\omega_{\text{ME}} = \sqrt{\omega_x \omega_y - \frac{\gamma B^2}{M_s C} (\omega_y \cos^2 \theta + \omega_x \sin^2 \theta)}. \quad (10)$$

Reflection R_{ω} , and transmission T_{ω} coefficients are then found via the well-known relations⁴⁵ as

$$R_{\omega} = \frac{(\eta_{\omega}^2 - 1) \cdot \sin(k_{\omega,x} \delta)}{(\eta_{\omega}^2 + 1) \cdot \sin(k_{\omega,x} \delta) + 2i \eta_{\omega} \cdot \cos(k_{\omega,x} \delta)}, \quad (11)$$

$$T_{\omega} = \frac{2i \eta_{\omega}}{(\eta_{\omega}^2 + 1) \cdot \sin(k_{\omega,x} \delta) + 2i \eta_{\omega} \cdot \cos(k_{\omega,x} \delta)}, \quad (12)$$

where δ is the thickness of the magnetic inclusion and $\eta_{\omega} = Z_{\omega, \text{ME}} / Z_0$.⁴⁶ In close proximity to the resonance, the impedance changes rapidly. For matching elastic properties of the magnetic and non-magnetic materials, Eq. (11) is expanded near ω_{ME} . In the limit $k_{\omega} \delta \ll 1$, we obtain

$$R_{\omega} = \frac{i \Gamma_{\text{R}} / 2}{(\omega - \omega_{\text{ME}}) + i \Gamma_{\text{R}} / 2} + R_0, \quad (13)$$

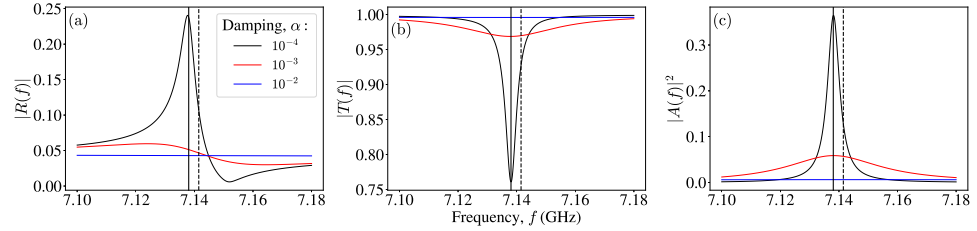


FIG. 2. The frequency dependence of the absolute values of (a) reflection and (b) transmission coefficients and (c) absorbance is shown for a 20nm thick magnetic inclusion. The vertical dashed and solid black lines represent the ferromagnetic resonance frequency ω_{FMR} and magneto-elastic resonance frequency ω_{ME} respectively. The non-magnetic and magnetic materials are assumed to be silicon nitride and cobalt, respectively, with parameters given in the text. The bias field is $\mu_0 H_B = 50\text{mT}$, which leads to $f_{\text{ME}} \approx 7.138\text{ GHz}$.

where R_0 represents the smooth non-resonant contribution due to elastic mismatch at the interfaces. In a system with no magnetic damping, the hybridization yields a resonance of finite linewidth Γ_R ,

$$\Gamma_R = \frac{\gamma B^2}{2M_s C} \sqrt{\frac{\rho}{C}} (\omega_x \cos^2 \theta + \omega_x \sin^2 \theta) \delta. \quad (14)$$

The origin of this linewidth can be explained as follows. Due to the magneto-elastic coupling incident propagating acoustic modes can be converted into localised magnon modes. These modes in turn either decay due to the Gilbert damping or are re-emitted as phonons. The rates of these transitions are proportional to Γ_{FMR} and Γ_R , respectively, and the total decay rate is $\Gamma = \Gamma_R + \Gamma_{\text{FMR}}$. This is similar to resonant scattering in quantum theory⁴⁷, such that Γ_R and Γ_{FMR} are analogous to the elastic (Γ_e) and inelastic (Γ_i) linewidths respectively. When $\alpha = 0$, Γ_{FMR} vanishes, and $\Gamma = \Gamma_R$.

Acoustic waves in the geometry of Fig. 1 can be scattered via several channels. E.g. in a non-magnetic system ($B = 0$), elastic mismatch can yield Fabry-Pérot resonance due to the quarter wavelength matching of δ and the acous-

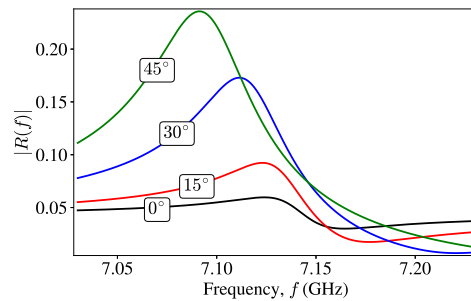


FIG. 3. Peak $R(f)$ is enhanced and slightly shifted to the left in the oblique incidence geometry ($\theta > 0^\circ$). Coloured curves represent specific incidence angles sweeping 0° to 45° . Moderate Gilbert damping of $\alpha = 10^{-3}$ is assumed.

tic wavelength. However, this occurs at very high frequencies, which we do not consider here. To understand the resonant magneto-elastic response, it is instructive to consider first the case of normal incidence ($\theta = 0$), when the demagnetising energy takes a simplified form due to the lack of immediate interfaces to form surface poles in y the direction, so that $N_x = 1$ and $N_y = 0$. Including magneto-elastic coupling ($B \neq 0$), we plot the frequency dependence of R_ω and T_ω using Eq. (11) and (12) in Fig.2. To gain a quantitative insight, we analysed a magnetic inclusion made of cobalt ($\rho = 8900\text{kgm}^{-3}$, $B = 10\text{MPa}$, $C = 80\text{GPa}$, $\gamma = 176\text{GHzT}^{-1}$, $M = 1\text{MAm}^{-1}$), embedded into a silicon nitride non-magnetic matrix ($\rho_0 = 3192\text{kgm}^{-3}$, $C_0 = 298\text{GPa}$). To highlight the resonant behaviour, we first suppress α to 10^{-4} . The reflection coefficient exhibits an asymmetric non-monotonic dependence, shown as a black curve in Fig.2(a), characteristic of Fano resonance.^{27,41} This line shape can be attributed to coupling between the discrete FMR mode of the magnetic inclusion and the continuum of propagating acoustic modes in the surrounding non-magnetic material.⁴¹ If the two materials had matching elastic properties, R_ω would exhibit a symmetric Breit-Wigner lineshape.⁴⁷ The transmission shown in Fig.2(b) exhibits an approximately symmetric dip near the resonance.⁴⁸ The absorbance $|A_\omega|^2 = 1 - |R_\omega|^2 - |T_\omega|^2$, shown in Fig.2(c) exhibits a symmetric peak, since the acoustic waves are damped in our model only due to the coupling with spin waves.

To consider how the magneto-elastic resonance is affected by the damping, we also plot the response for α of 10^{-3} and 10^{-2} , red and blue curves in Fig.2, respectively. An increase of α from 10^{-4} to 10^{-3} significantly suppresses and broadens the resonant peak. For a more common, realistic value of 10^{-2} the resonance is quenched entirely. A stronger magneto-elastic coupling (i.e. high values of B) could, in principle, countermand this suppression. This, however, is also likely to enhance the phonon contribution to the magnetic damping, leading to a correlation between B and α observed in realistic magnetic materials.⁴⁹

To characterise the strength of the Fano resonance, we note that the fate of the magnon excited by the incident acoustic wave is decided by the relation between the emission rate Γ_R ,

see Eq. (14), and absorption rate Γ_{FMR} . Hence, we introduce the respective figure of merit as $\Upsilon = \Gamma_{\text{R}}/\Gamma_{\text{FMR}}$. This quantity depends upon the material parameters, device geometry, and bias field. As seen from the first terms on the l.h.s. of Eqs. (6) and (7), the relation between the dynamic magnetisation components $m_{x,y}$ are determined by the quantities ω_x and ω_y . Equating these terms, one finds $m_x \propto m_y \sqrt{\omega_y/\omega_x}$, i.e. the precession of \mathbf{m} is highly elliptical,⁵⁰ due to the demagnetising field along x . This negatively affects the phonon-magnon coupling for normal incidence ($k_y = 0$): the acoustic wave couples only to m_x , as given by the second term in Eqs. (6) and (7). One way to mitigate this is to increase H_B , moving the ratio ω_y/ω_x closer to 1 and thus improving the figure of merit. To compare different magneto-elastic materials, the dependence on the layer thickness δ and elastic properties of the non-magnetic matrix (i.e. ρ_0 and C_0) can be eliminated by calculating a ratio of the figures of merit for the compared materials. The comparison can be performed either at the same value of the bias field, or at the same operating frequency. The latter situation is more appropriate for a device application, but to avoid unphysical parameters, we present our results for the same $\mu_0 H_B$. An example of such comparisons for yttrium iron garnet (YIG), cobalt (Co) and permalloy (Py) is offered in Table I.

TABLE I. Comparison of the figure of merit Υ for different materials, assuming $\delta = 20\text{nm}$, $\mu_0 H_B = 50\text{mT}$ and a silicon nitride non-magnetic matrix.

Parameters	YIG	Co	Py
$\Upsilon(\theta = 0^\circ)$	4.3×10^{-3}	1.5×10^{-3}	1.9×10^{-4}
Γ_{R} (GHz)	3.0×10^{-6}	6.5×10^{-3}	1.4×10^{-4}
Γ_{FMR} (GHz)	7.0×10^{-4}	4.3	0.74
$\Upsilon(\theta = 30^\circ)$	8.1×10^{-3}	1.1×10^{-2}	1.1×10^{-3}
Γ_{R} (GHz)	5.6×10^{-6}	4.7×10^{-2}	8.0×10^{-4}
Γ_{FMR} (GHz)	7.0×10^{-4}	4.3	0.74
ω_{ME} (GHz)	0.48	7.14	6.26
B (MJm^{-3})	0.55	10	-0.9
E (GPa)	74	80	50
ρ (kgm^{-3})	5170	8900	8720
α	0.9×10^{-4}	1.8×10^{-2}	4.0×10^{-3}
M_s (kAm^{-1})	140	1000	760

Another way to improve Υ is to employ the oblique incidence ($\theta \neq 0$), in which the acoustic mode is also coupled to the magnetisation component m_y . The latter is not suppressed by the demagnetisation effects if $N_y \ll 1$. The resulting enhancement in Υ is reflected in the full equation by the inclusion of ω_x and ω_y from Γ_{R} ,

$$\Upsilon = \frac{\Gamma_{\text{R}}}{\Gamma_{\text{FMR}}} = \frac{\gamma \delta}{2} \sqrt{\frac{\rho_0 B^2 (H_B \cos^2 \theta + M_s \sin^2 \theta)}{C_0 \alpha C M_s^2}}, \quad (15)$$

where $\omega_x \gg \omega_y$ and $H_B \ll M_s$ is assumed. For small θ , the approximation $N_x \simeq 1$ and $N_y \simeq 0$ still holds. As a result, non-zero θ increases peak reflectivity, as seen in Fig.3. Note also the θ dependence of the resonant frequency ω_{ME} as reflected in Eq. (10). Though larger incidence may be hard to imple-

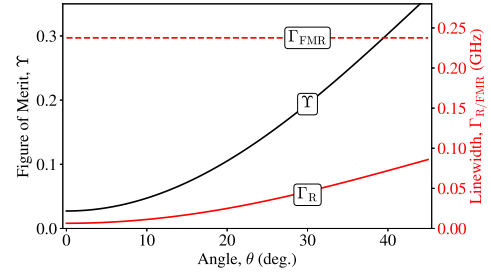


FIG. 4. Figure of merit Υ , and radiative linewidth Γ_{R} , are both enhanced in the oblique incidence geometry ($\theta > 0^\circ$). Ferromagnetic linewidth Γ_{FMR} remains unchanged. Co is assumed with $\alpha = 10^{-3}$.

ment in a practical device, resonant scattering is still significantly enhanced at smaller angles: $\theta \approx \sqrt{H_B/M_s} \ll 1$ yields a twofold increase in Γ_{R} and Υ , as illustrated in Fig.4.

Above, we have focused on the simplest geometry that admits full analytic treatment. To implement our idea experimentally, particular care should be taken about the acoustic waves polarization and propagation direction relative to the direction of the magnetization. Indeed, our choice maximises magnetoelastic response. If however, the polarization is orthogonal to the bias field H_B , i.e. $U_z = 0$, the coupling would be second-order in magnetization components $m_{x,y}$, and would not contribute to the linearized LLG equation. Furthermore, we have neglected the exchange and magneto-dipolar fields that could arise due to the non-uniformity of the magnetization. To assess the accuracy of this approximation, we note that the length scale of this non-uniformity is set by the acoustic wavelength λ , of about 420nm for our parameters rather than by the magnetic slab thickness δ . The associated exchange field is $\mu_0 M_s (k l_{\text{ex}})^2 \simeq 9\text{mT}$. The k -dependent contributions to the magneto-dipole field vanish at normal incidence but may become significant at oblique incidence, giving $\mu_0 M_s k_y \delta \simeq 98\text{mT}$ at $\theta = 15^\circ$. In principle, these could increase the resonant frequency of the slab by a few GHz but would complicate the theory significantly. The detailed analysis of the associated effects is beyond the scope of this report.

In summary, we have demonstrated that the coupling between the magnetisation and strain fields can be used to control acoustic waves by magnetic inclusions. We show that the frequency dependence of the waves' reflection coefficient from the inclusions has a Fano-like lineshape, which is particularly sensitive to the magnetic damping. Figure of merit is introduced to compare magnetoelastic materials and to characterize device performance. In particular, the figure of merit is significantly enhanced for oblique incidence of acoustic waves, which enhances their coupling to the magnetic modes. We envision that further routes may be taken to transform our prototype designs into working devices, such as forming a magneto-acoustic metamaterial to take advantage of spatial resonance.

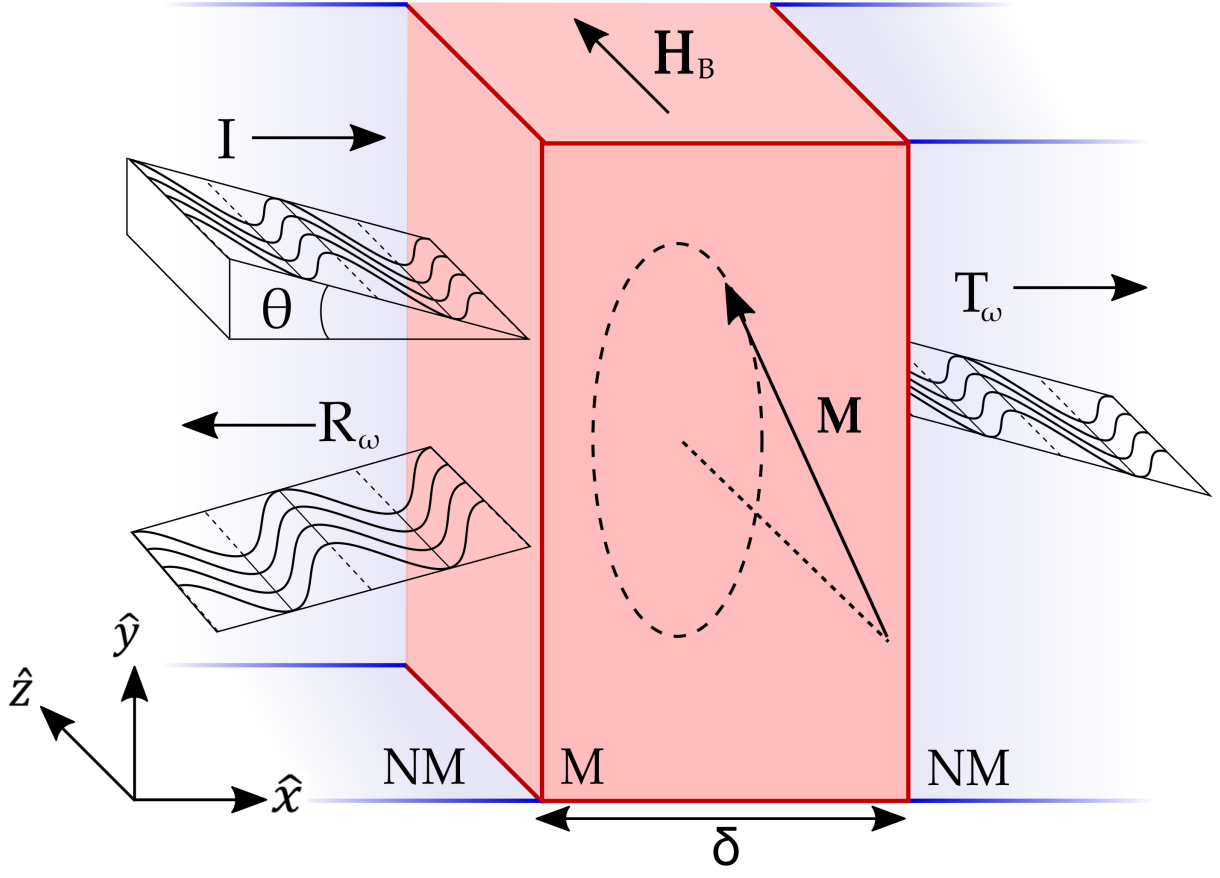
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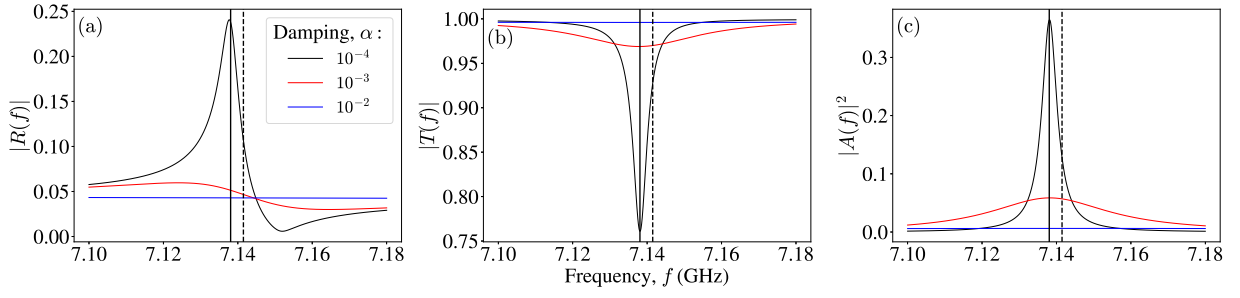
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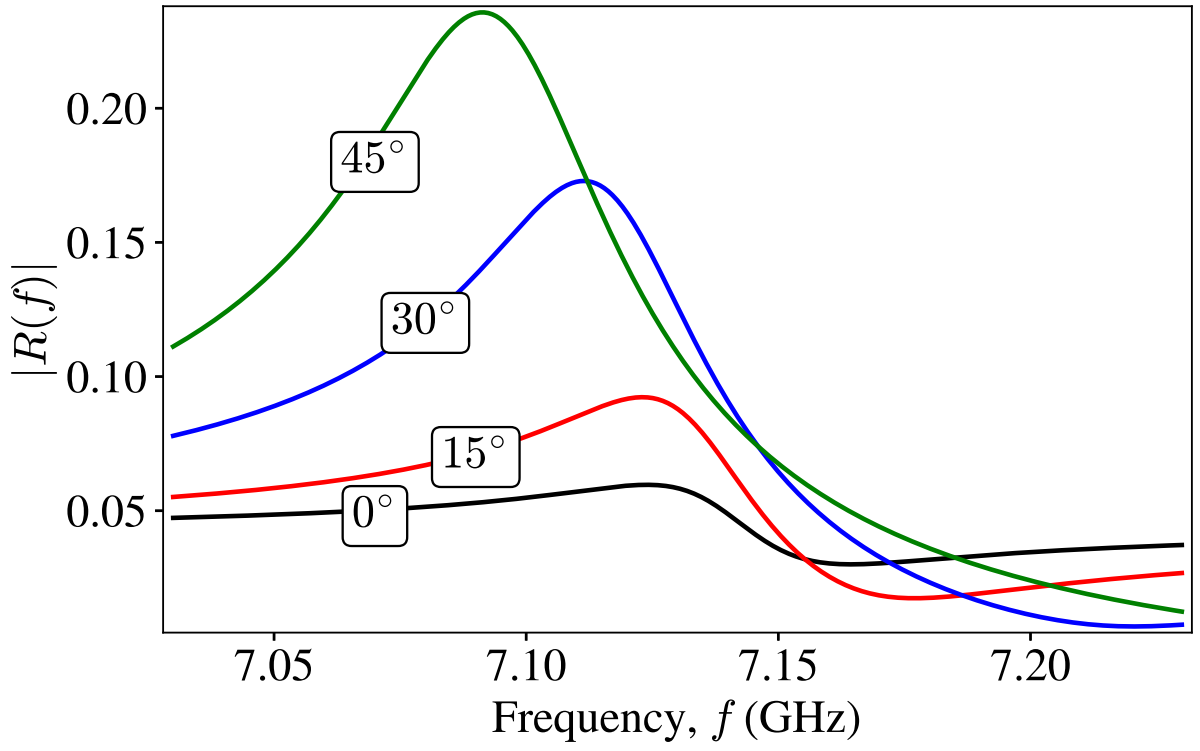
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