Solid State Communications 198 (2014) 35–39

Contents lists available at ScienceDirect

Solid State Communications

journal homepage: www.elsevier.com/locate/ssc

Actuation, propagation, and detection of transverse magnetoelastic waves in ferromagnets

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ARTICLE INFO

Article history:
Received 27 June 2013
Received in revised form
21 August 2013
Accepted 7 October 2013
by Sebastian T.B. Goennenwein
Available online 15 October 2013

Keywords:
A. Ferromagnets
D. Ultrasound
D. Magnetoelastic coupling
D. Spin pumping

ABSTRACT

We study propagation of ultrasonic waves through a ferromagnetic medium with special attention to the boundary conditions at the interface with an ultrasonic actuator. In analogy to charge and spin transport in conductors, we formulate the energy transport through the system as a scattering problem. We find that the magneto-elastic coupling leads to a non-vanishing magnetic (elastic) energy accompanying the acoustic (spin) waves with a resonantly enhanced effect around the anti-crossing in the dispersion relations. We demonstrate the physics of exciting magnetization dynamics via acoustic waves injected around the ferromagnetic resonance frequency.

1. Introduction

While the exchange interaction is the largest energy scale of ferromagnets, explaining Curie temperatures of up to 1000 K, the static equilibrium and dynamic properties of the magnetization field in ferromagnetic materials are governed by the dipolar and crystal anisotropy fields [1]. Since the total angular momentum of an isolated system is conserved, any change of the magnetization exerts a torque on the underlying lattice, as measured by Einstein and de Haas [2]. Vice versa, a rotating lattice can magnetize a demagnetized ferromagnet [3]. The coupled equations of motion of lattice and magnetization fields have been treated in a seminal paper reported by Kittel [4]. The magnetoelastic coupling parameters are material constants well known for many ferromagnets [5].

Interest in magnetoelastic coupling has recently been revived in the context of the "spin mechanics" concept covered by the present special issue [Editorial SSC]. Here we are interested in the magnetization dynamics acoustically induced by injecting ultrasound into ferromagnets by piezoelectric actuators as bulk [6] or surface [7] plane acoustic waves. The magnetization dynamics in these experiments is conveniently detected by spin pumping [8] into a normal metal that generates a voltage signal via the inverse spin Hall effect [9].

Some of the consequences of magnetoelastic coupling have already been investigated theoretically [4,10,11] and experimentally [12,13] in the literature. While the coupled magnetoelastic dynamics has been well understood decades ago [4,10,14], much less attention has been devoted to the interfaces that are essential in order to understand modern experiments on nanostructures and ultrathin films. The Landauer–Büttiker electron transport formalism based on scattering theory is well suited to handle these issues thereby helping to understand many problems in mesoscopic quantum transport and spintronics [15,16]. Here we formulate scattering theory of lattice and magnetization waves in ferromagnets with significant magnetoelastic coupling. Rather than attempting to describe concrete experiments, we wish to illustrate here the usefulness of this formalism for angular momentum and energy transport.

We consider magnetization dynamics actuated by ultrasound for the simplest possible configuration in which the magnetization direction is parallel to the wave vector of sound with transverse polarization (shear waves). The corresponding bulk propagation of magnetoelastic waves was treated long ago by Kittel [4] who demonstrated that the axial symmetry reduces the problem to a quadratic equation. The injected acoustic energy is partially transformed into magnetic energy by the magnetoelastic coupling that can be detected by spin pumping into a thin Pt layer. For this symmetric configuration and to leading order a purely AC voltage is induced by the inverse spin Hall effect (ISHE) [17] that might be easier to observe by acoustically induced rather than rf radiation induced spin pumping, since in the former the Pt layer is not

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0038-1098/$ - see front matter
http://dx.doi.org/10.1016/j.ssc.2013.10.007

directly subjected to electromagnetic radiation. The configuration considered by Uchida et al. [6], in which pressure waves generate a DC ISHE voltage by a magnetization parallel to the interfaces, will be discussed elsewhere.

2. Kittel's equations

We consider a ferromagnet with magnetization texture $\mathbf{M}(\mathbf{r},t)$ with constant saturation magnetization $|\mathbf{M}| = M_0$. In the following we consider small fluctuations around the equilibrium magnetization $M_0\mathbf{z}$. The classical Hamiltonian can be written as the sum of different energies

$$\mathcal{H} = \mathcal{H}_Z + \mathcal{H}_{ex} + \mathcal{H}_{me} + \mathcal{H}_p \tag{1}$$

The magnetic Zeeman energy reads

$$\mathcal{H}_Z = \frac{\omega_0}{2\gamma M_0}(M_x^2 + M_y^2) \tag{2}$$

where $\gamma = |\gamma|$ is the gyromagnetic ratio, $\omega_0 = \mu_0 \gamma H$ is the magnetic resonance frequency for an effective magnetic field $H\mathbf{z}$ and μ_0 the permeability of free space. The exchange energy cost of the fluctuations reads

$$\mathcal{H}_{ex} = \frac{A}{M_0^2}[(\nabla M_x)^2 + (\nabla M_y)^2] \tag{3}$$

where A is the exchange constant. The magnetoelastic energy for cubic crystals and magnetization in the z-direction can be parameterized by the magnetoelastic coupling constant b_2

$$\mathcal{H}_{me} = \frac{b_2}{M_0}\left(M_x \frac{\partial R_x}{\partial z} + M_y \frac{\partial R_y}{\partial z}\right) \tag{4}$$

where $\mathbf{R} = (R_x, R_y, 0)$ is the displacement vector of a transverse lattice wave propagating in the z-direction. $\mathcal{H}_{me}$ can be interpreted as a Zeeman energy associated with a dynamic transverse magnetic field $b_2 \partial_z \mathbf{R}$. The corresponding elastic energy reads

$$\mathcal{H}_p = \frac{\rho}{2}\dot{\mathbf{R}}^2 + \frac{\alpha}{2}\left[\left(\frac{\partial R_x}{\partial z}\right)^2 + \left(\frac{\partial R_y}{\partial z}\right)^2\right] \tag{5}$$

in terms of the mass density ρ and shear elastic constant α.

The total Hamiltonian $\mathcal{H}$ defines the equations of motion of the coupled $\mathbf{R}$ and $\mathbf{M}$ fields. The results in momentum and frequency space $X(t) = x(k,\omega)e^{i(kx-\omega t)}$ can be simplified by introducing circularly polarized phonon and magnon waves $m^{\pm} = m_x + i\sigma m_y$, $r^{\pm} = r_x + i\sigma r_y$ ($\sigma = \pm 1$), leading to [4]

$$\begin{pmatrix} i(\omega - \sigma\omega_m) & \sigma\gamma b_2 k \\ i\frac{b_2 k}{M_0} & \omega^2\rho - k^2\alpha \end{pmatrix}\begin{pmatrix} m^{\sigma} \\ r^{\sigma} \end{pmatrix} = 0 \tag{6}$$

where $\omega_m = \omega_0 + Dk^2$ and $D = 2A\gamma/M_0$ is the spin wave stiffness. This secular equation is quadratic in k^2 with 4 roots ($s = \pm 1$):

$$(k_s^{\sigma})^2 = \frac{\rho\omega^2}{2\alpha} - \frac{\omega_0 - \sigma\omega}{2D} + \frac{\gamma b_2^2}{2\alpha M_0} + s\sqrt{\Delta^{\sigma}} \tag{7}$$

with discriminants

$$\Delta^{\sigma} = \left(\frac{\rho\omega^2}{2\alpha} - \frac{\omega_0 - \sigma\omega}{2D} + \frac{\gamma b_2^2}{2\alpha D M_0}\right)^2 + \frac{\omega^2\rho}{\alpha D}(\omega_0 - \sigma\omega). \tag{8}$$

The corresponding eigenstates are given by the spinor

$$\psi_s^{\sigma}(\omega) = \begin{pmatrix} m_s^{\sigma} \\ r_s^{\sigma} \end{pmatrix} = N_s^{\sigma}\begin{pmatrix} M_0 \\ ib_2 k_s^{\sigma}/(\omega^2\rho - (k_s^{\sigma})^2\alpha) \end{pmatrix}, \tag{9}$$

where $N_s^{\pm}$ is a dimensionless normalization factor.

The dispersion is plotted in Fig. 1 for the parameters appropriate for Yttrium Iron Garnet (YIG): $M_0 = 1.4\times 10^5$ A/m, $b_2 = 5.5\times 10^5$ J/m^3, $H = 8\times 10^4$ A/m, $D = 8.2\times 10^{-6}$ m^2/s, $\gamma = 2.8\times 10^{10}$ Hz T^{-1}, $\rho = 5170$ kg/m^3, $\alpha = 7.4\times 10^{10}$ Pa [18–20] and $\mu_0 = 4\pi\times 10^{-7}$ NA^{-2}. In Fig. 1(a) we plot the solutions for waves rotating with the magnetization that appear to be completely phonon (small dispersion) or magnon like (large dispersion). The latter are evanescent ($k^2 < 0$) below the spin wave gap ω_0. The low-frequency anticrossing is better seen in Fig. 1(c) in which the momentum is plotted on an expanded scale. Spin waves precessing against the magnetization, $\sigma = -1$, are always evanescent and there is no (anti)crossing with the propagating phonons. When the magnetoelastic coupling is switched off ($b_2 \to 0$) and $\alpha > 4\rho D\omega_0$ the pure lattice wave $\omega_p = \sqrt{\alpha/\rho}\,k$ and spin wave $\omega_m = \omega_0 + Dk^2$ dispersions may cross twice

$$\omega_c = \frac{\alpha}{2\rho D}\left(1 \pm \sqrt{1 - \frac{4\rho D\omega_0}{\alpha}}\right) \tag{10}$$

$$\overset{4\rho D\omega_0 \ll \alpha}{=} \begin{cases} \omega_0 \\ \dfrac{\alpha}{\rho D} \end{cases} = \begin{matrix} 2.8\text{ GHz} \\ 1.7\text{ THz} \end{matrix} \tag{11}$$

where in the second step we take the limit of small D. Around these degeneracy points, of which only the low frequency one is relevant here, the effects of the magnon–phonon coupling are most pronounced. In the zero frequency limit the solution with $(k_s^{\sigma})^2 \to 0$ represents a phonon mode with zero wave number. $(k_s^{\sigma})^{\pm} = \gamma b_2^2/(\alpha M_0) - \omega_0/D$ is a purely evanescent magnon for small b_2 that in principle may become a real excitation when the coupling of the lattice is strong enough to overcome the spin wave gap.

3. Energy flux

Energy conservation implies $\nabla\cdot\vec{F} = -\partial\mathcal{H}/\partial t$, where the energy flux $\vec{F} = F\hat{\mathbf{z}}$ consists of phonon and magnon contributions. In time and position space [21]

$$F(z,t) = -\int dz \frac{\partial\mathcal{H}}{\partial t} = -\frac{2A}{M_0^2}\frac{\partial M_x}{\partial z}\frac{\partial M_x}{\partial t} - \left(\alpha\frac{\partial R_x}{\partial z} + \frac{b_2}{M_0}M_x\right)\frac{\partial R_x}{\partial t} + (x \longleftrightarrow y). \tag{12}$$

For a plane wave oscillating with frequency ω

$$X(z,t) = x(z,\omega)e^{-i\omega t} + x^*(z,\omega)e^{i\omega t} \tag{13}$$

$$X(z,t) = x(\omega)e^{i(kx-\omega t)} + x^*(\omega)e^{-i(kx-\omega t)} \tag{14}$$

the time-averaged energy flux reads

$$\overline{F}(z)_x = -2\omega\,\mathrm{Im}\left[\frac{D}{\gamma M_0}m_x\partial_z m_x^* + \alpha r_x \partial_z r_x^* + \frac{b_2}{M_0}r_x m_x^*\right]. \tag{15}$$

In the absence of magnetoelastic coupling, pure phonon and magnon waves

$$\psi_s^{(m)} = N_s^{(m)} M_0 \begin{pmatrix} 1 \\ 0 \end{pmatrix} e^{i(k_s z - \omega t)} \tag{16}$$

$$\psi_0^{(p)} = N_p \begin{pmatrix} 0 \\ k^{-1} \end{pmatrix} e^{i(kz-\omega t)} \tag{17}$$

carry, respectively, the energy fluxes $\overline{F}^{(p)}$ and $\overline{F}_s^{(m)}$:

$$\overline{F}^{(p)} = N_p^2 2\alpha\omega/k = N_p^2 2\alpha\sqrt{\alpha/\rho} \tag{18}$$

$$\overline{F}_s^{(m)} = (N_s^{(m)})^2 \frac{2DM_0}{\gamma}\omega k_s \tag{19}$$

$$\overline{F}_s^{(m)} = (N_s^{(m)})^2 \begin{cases} \dfrac{2D}{\gamma M_0}\omega\sqrt{|s\omega - \omega_0|/D} & \text{for } s\omega > \omega_0, \\ 0 & s\omega < \omega_0. \end{cases} \tag{20}$$