Spin-wave excitations in multiferroic heterostructures and CoFeB/YIG bilayers

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A doctoral dissertation completed for the degree of Doctor of Science (Technology) to be defended, with the permission of the Aalto University School of Science, at a public examination held at the lecture hall U1 of the Main Building (Otakaari 1) on 12 October 2018 at 12.

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Spin waves are collective excitations in magnetic materials that can exhibit long coherence lengths, high group velocities, and wavelengths down to the nanometer scale. The properties of spin waves could be exploited in low-power wave-like computing and in other devices such as microwave filters and data storage elements. In contrast to the motion of electrons, spin waves propagate via transmission of angular momentum. Power dissipation by Joule heating can thus be avoided.

In this thesis, I present results on spin-wave excitation and confinement in strain-coupled multiferroic heterostructures and YIG/CoFeB bilayers. The multiferroic heterostructures consist of a ferroelectric \( BaTiO_3 \) substrate with a regular pattern of ferroelastic stripe domains and a ferromagnetic film. Strain coupling at the interface imprints an anisotropy pattern in the ferromagnetic film via inverse magnetostriction. As a result, the domain pattern of the \( BaTiO_3 \) substrate is fully transferred to the ferromagnetic film and magnetic domain walls are firmly pinned by the induced anisotropy boundaries. The pinned magnetic domain walls can be used to excite short-wavelength spin waves using an applied electric current. In this case, the spin-polarized current drives the domain wall into oscillatory motion, causing spin-wave emission at the same frequency as the actuation signal. If the same multiferroic system is saturated by an external magnetic field and excited by a uniform microwave field, standing spin waves form within the domains and propagating spin wave are emitted from the anisotropy boundaries. Here, rotations of the anisotropy axis with respect to the direction of magnetization create a regular modulation of the effective field. The launching of spin waves from a single anisotropy boundary is explained by coupling between two forced magnetization precessions in neighboring stripe domains. A similar mechanism is also responsible for the excitation of perpendicular standing spin waves in exchange-coupled YIG/CoFeB bilayers. In this system, forced precession in YIG and CoFeB generate a dynamic exchange torque at the interface. Lastly, multiferroic heterostructures are used to study spin wave transmission through narrow and wide \( 90^\circ \) domain walls. Deterministic switching between these wall types is realized by realigning the magnetization from a head-to-tail (narrow wall) to a head-to-head (broad wall) configuration. Brillouin light scattering experiments and micromagnetic simulations demonstrate that this reversible switching effect toggles the transmission of spin waves from nearly 0% to 100%. Reconfigurable modulations of spin-wave transmission could be used in logic devices or non-volatile memory cells.
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This thesis consists of an overview and of the following publications which are referred to in the text by their Roman numerals.


VI Pavel Baláž, Sampo J. Hämäläinen & Sebastiaan van Dijken. Static properties and current-induced dynamics of pinned 90 magnetic domain walls under


Author’s Contribution

**Publication I: “Size Dependence of Domain Pattern Transfer in Multiferroic Heterostructures”**

The author contributed to data analysis and discussed the results with co-authors.

**Publication II: “Reversible Electric-Field-Driven Magnetic Domain-Wall Motion”**

The author developed the setup for controlled application of electric-field pulses inside the magneto-optical Kerr microscope. He wrote the script for data analysis and discussed the results with co-authors.

**Publication III: “Influence of elastically pinned magnetic domain walls on magnetization reversal in multiferroic heterostructures”**

The author developed the script for data analysis and discussed the results with co-authors.

**Publication IV: “Tunable short-wavelength spin wave excitation from pinned magnetic domain walls”**

The author modeled the dispersion relations of emitted spin waves and discussed the results with co-authors.
Author's Contribution


The author performed and analyzed the micromagnetic simulations and contributed to the broadband spin-wave spectroscopy experiments. He wrote the first draft of the manuscript.

Publication VI: “Static properties and current-induced dynamics of pinned 90 magnetic domain walls under applied fields: An analytic approach”

The author performed and analyzed the micromagnetic simulations.

Publication VII: “Programmable control of spin-wave transmission in a domain-wall spin valve”

The author performed the micromagnetic simulations, prepared the sample and contributed to the Brillouin light scattering experiments. He wrote the first draft of the manuscript.

Publication VIII: “Exchange-torque-induced excitation of perpendicular standing spin waves in nanometer-thick YIG films”

The author contributed to the micromagnetic simulations and the establishment of the spin-wave measurement setup. He discussed the results with co-authors.
Other publications by the author that are not included in the thesis.


Other publications by the author
1. Introduction

Digital information and communication technologies are an important part of our daily lives. The rapid growth rate in computational capabilities was forecasted by Gordon Moore in 1965 who predicted that the transistor count in integrated circuits would double every two years [80]. This exponential growth is progressing even during present days. However, it has been estimated that this progress will slow down in the near future as current commercial transistor dimensions reached the 14 nm node. Besides lithography, challenges involve quantum mechanical tunneling through the oxide layer of the transistor and high Joule losses. Spin waves offer an alternative approach to computation, including parallel processing. Since spin waves propagate by transfer of angular momentum instead of electron motion, power dissipation by Joule heating can be avoided.

In spin wave or so called magnonic devices, the information is carried by the amplitude or phase of propagating waves. One proposed transistor-like device that can be used as logic gate relies on constructive or destructive interference of spin waves in a Mach-Zehnder interferometer [95]. In this system, a spin wave is divided into two branches and merged at the output port. By controlling the phase of the spin wave in one arm, constructive or destructive interference can be attained at the output. The phase-shift can be induced, for example, by formation of a domain wall [42] or the use of Oersted fields [95]. Alternatively, spin wave transport can be controlled in magnonic crystals. In magnonic crystals, magnetic properties are modulated periodically, causing the opening of bandgaps when the wave-length matches the period of the modulation crystal [58]. Active manipulation of spin waves in magnonic crystals has been realized using, for example, electric currents [18], optical pulses [112], or magnetic switching in bicomponent lattices [104].

In this thesis, I present a new structure for active and efficient manipulation of propagating spin waves. The structure consists of two parallel pinned magnetic domain walls and is experimentally realized in a strain-coupled BaTiO3/CoFeB multiferroic bilayer with regular rotations of uniaxial magnetic anisotropy. The abrupt rotations of uniaxial anisotropy strongly pin the magnetic domain walls on top of perfectly straight ferroelectric domain the boundaries. At zero magnetic...
This page contains a detailed introduction to the topic of magnetic spin waves, focusing on the excitation and detection of short-wavelength spin waves at the nanoscale. The text explains how magnetization rotation within the domain walls is fixed at 90°. However, the width of the domain walls can be deterministically and reversibly switched from narrow to wide by realigning the magnetization from a head-to-tail to a head-to-head or tail-to-tail configuration. Broad magnetic domain walls are fully transparent to propagating spin waves while the transmission of spin waves is reduced to nearly zero after the magnetization configuration is switched to narrow domain walls. The system with two parallel pinned domain walls, which is easily switched by magnetization reversal of the central domain, could be exploited in magnonic logic devices or as a non-volatile memory cell. In analogy to the magnetic spin valve structure for electrons, I refer to this structure as magnetic spin wave valve.

Besides active manipulation of spin waves, efficient excitation and detection of short-wavelength spin waves is essential for the realization of magnonic devices at the nanoscale. Sub-100 nm spin waves have been excited by uniform microwave magnetic fields using grating coupler structures [118, 119, 76]. Short wavelength spin waves can also be attained by exploiting tapered [28] or thickness modulated waveguides [100], isolated edges or antidots [27], or microwave to spin wave transducers [4].

In this thesis, I present several approaches to excite short wavelength spin waves. The first approach utilizes a pinned 90° magnetic domain wall in an anisotropy-modulated BaTiO3/CoFe multiferroic heterostructure. When a spin-polarized current is applied perpendicular to the domain wall, it is forced to oscillate at the same frequency because of a spin-transfer torque effect. The back-and-forth motion of the magnetic domain wall results in the excitation of propagating spin waves. Efficient spin-wave excitation can be sustained up to higher frequencies if the ferromagnetic film is patterned into nanowires. In this case, higher-order domain-wall resonances are excited, enabling downscaling of the spin-wave wavelength to about 20 nm at high frequencies. The multiferroic heterostructure is also exploited to excite and confine spin waves in magnetic saturation, i.e., without any domain walls being present. Under this condition, rotations of the magnetic anisotropy axis create a regular modulation of the effective magnetic field. In anisotropy domains with smaller effective magnetic field, higher-order standing spin waves are exited by uniform microwave fields. Since no propagating modes exist in neighboring domains with larger effective field, these modes are fully confined. In contrast, when standing spin wave modes are exited in domains with higher effective field, propagating spin waves are launched into the neighbouring domains. The wavelength of these propagating spin waves can be tuned down to about 100 nm by rotation of the external magnetic bias field. Furthermore, I demonstrate that uniform microwave fields cause the emission of spin waves from a single magnetic anisotropy boundary. This excitation mechanism is explained by coupling between the forced magnetization precessions in two neighbouring domains. A similar principle also results in an efficient excitation of perpendicular standing spin waves in exchange coupled YIG/CoFeB bilayers. In this case, the large difference in saturation
magnetization between the two layers forces their magnetization to precess with different amplitudes. As a result, a dynamic exchange torque is generated at the interface, causing the emission of spin waves. If the wavelength of the waves matches one of the confinement conditions, a perpendicular standing spin wave is formed. Efficient excitation of short-wavelength spin waves in YIG at elevated frequencies, as demonstrated, is relevant for the development of nanoscale magnonic devices.
2. Ferromagnetism

Ferromagnetic materials are characterized by parallel alignment of neighbouring magnetic moments which leads to spontaneous magnetization. The magnetic moment arises from the intrinsic spin and orbital angular momentum of electrons. By introducing coupling between neighbouring moments, magnetization can retain a finite value even at zero magnetic field. The formation of spontaneous magnetization is induced by exchange interactions, which in ferromagnetic materials prefer parallel alignment of neighbouring spins. When a ferromagnetic material is heated above the Curie temperature, thermal excitations dominate over the exchange interactions. In this paramagnetic state, the random fluctuations of the spins generate a disordered state with no net magnetic moment.

Studies of magnetically ordered materials can be divided into four general approaches: atomic-level theory (~1 nm), micromagnetics (~1-1000 nm), domain theory (~1 μm-1 mm) and phase theory (0.1 mm) [44]. In the first case, discreet magnetic moments $m$ are used to describe the origin of magnetism, calculate anisotropy, and evaluate thermodynamics fluctuations. However, when the system size increases, the number of interactions in many-body systems grows exponentially and atomic-level theory becomes challenging. To overcome this issue, micromagnetic theory omits the atomistic structure and approximates the system by continuum theory. Here, the spin distribution is described by a continuous variable $M(r)$ as a function of position $r$. This tool reliably resolves the properties and internal structure of domains and domain walls. However, micromagnetism is insufficient when studying larger samples on the millimeter scale. In this case, the inner magnetic microstructure cannot be experimentally probed [44]. To resolve details of the domain wall structure and interconnection between domain walls, micromagnetic theory is combined with discreet and uniformly magnetized domains in domain theory [24, 53]. The last approach, phase theory, omits the details of the domain texture and concentrates on the volume distribution of an ensemble of domains. This is particularly useful when studying large soft samples where the texture and energies of individual domain walls are unimportant.
2.1 Continuum theory and micromagnetism

In this thesis, the described magnetic systems have micrometer length scales and hence their properties are suitably described by micromagnetic theory. To solve static magnetic properties, the energy terms arising from the external magnetic field, dipole-dipole interactions, exchange coupling and anisotropy are calculated and the total energy is minimized to find the local ground state. The minimal energy state can be simulated in two ways either by calculating the total energy of the system or by extracting the effective field that arises from the energy terms [79]. In the former approach, the total energy of the system is calculated and minimized by using, for example, a conjugate-gradient method [31]. This technique is useful in finding the static equilibrium state of a magnetic system [79], but it is not suitable for dynamic simulations. An alternative approach relies on the calculation of an effective field. This effective field can be directly calculated from the functional derivative of the energy density as

\[ B_{\text{eff}} = -\frac{1}{M_s} \frac{\delta e}{\delta m} \]  (2.1)

where \( M_s \) is the saturation magnetization and \( m \) is the normalized magnetic moment. In the field-approach, the effective field is evaluated directly from the magnetization distribution and it does not require the total energy of the system to be calculated [79]. In this case, the static equilibrium state is solved by minimizing the magnetic torque applied to each magnetic moment

\[ \frac{dm}{dt} = \bar{\tau} = -\gamma m \times B_{\text{eff}} \]  (2.2)

where \( \gamma \) is the gyromagnetic ratio. If the torque is non-zero, the magnetic moment moves towards the effective field. This implies that the energy of the system is not minimized.

2.2 Exchange interaction

Exchange interaction is a short range term that prefers (anti-)parallel orientation of neighbouring spins. The Heisenberg Hamiltonian describes the exchange interaction between pairs of spins \( S_i, S_j \)

\[ H = -\sum_{i,j} J_{i,j} <S_i \cdot S_j> \]  (2.3)

where \( J \) is the exchange constant. Positive values of \( J \) represent ferromagnetic ordering where neighbouring spins are aligned parallel, and negative values of \( J \) result in antiferromagnetic order. This discrete model can be approximated using Taylor series and a continuous moment distribution. The effective field can be calculated from the variational derivative as [44]

\[ B_{\text{ex}} = 2 \frac{A_{\text{ex}}}{M_{\text{sat}}} (\nabla m(r))^2 \]  (2.4)
where $M_{sat}$, $A_{ex}$ and $m$ are the saturation magnetization, exchange constant and unit vector of magnetization, respectively.

2.3 Zeeman energy

Zeeman energy describes the interaction between an external magnetic field $H_{ext}$ and magnetization. In ferromagnetic materials, this interaction prefers parallel alignment between the moments and external field $H_{eff}$ and the corresponding energy density is

$$e_z = -\mu M_{sat}(m \cdot H_{ext})$$

where $\mu$ is vacuum permeability.

2.4 Magnetostatic energy

Each magnetic moment in a sample induces a dipolar field that interacts with other moments in the system. Dipolar interactions are effective over longer ranges than exchange interactions. As each moment in the system interacts with other moments, this energy term is computationally the most demanding task in micromagnetic simulations [108]. The magnetostatic field arising from magnetic moments within sample volume $V$ can be calculated as [108]

$$B_{ms}(r) = \frac{\mu_0}{4\pi} \int_V \nabla \nabla \cdot M_s(r) \cdot m(r') \, dr'. \quad (2.6)$$

2.5 Anisotropy energy

Magnetic anisotropy energy arises from spin-orbit coupling which induces preferential cones, axes or planes for the direction of magnetization. The existence of favorable magnetization directions is a prerequisite for the formation of magnetic domains. Magnetic anisotropy can have different origins, including crystal symmetry, mechanical strain, or the shape of patterned structures.

2.5.1 Shape anisotropy

The dipole-dipole interaction between magnetic moments introduces a shape dependant anisotropy in a ferromagnetic sample. The magnetization aligns to minimize the magnetostatic energy over the sample volume. For example, the magnetic moments of a narrow nanowire align parallel to the wire to minimize the demagnetization field. More generally, the moments tend to align parallel to the edges of a patterned magnetic structure. In most cases, shape anisotropy is
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non-analytical. The effect can only be solved analytically for highly symmetrical or infinite shapes, e.g., infinite thin films or ellipsoids.

2.5.2 Magnetocrystalline anisotropy

Magnetocrystalline anisotropy arises from spin-orbit coupling in the lattice of a ferromagnet. It is often strong in single crystal magnetic materials. For cubic lattices, such as Fe, the magnetocrystalline anisotropy energy density can be written as [44]

\[ e = K_{c1}(m_1^2m_2^2 + m_1^2m_3^2 + m_2^2m_3^2) + K_{c2}m_1^2m_2^2m_3^2 \] (2.7)

where \( m_i \) are the normalized magnetization components along the cubic axes (\( m_i = M_i/M_s \)). The signs of the anisotropy constants \( K_{c1} \) and \( K_{c2} \) define the easy axes of magnetization in the crystal.

2.5.3 Magnetoelastic anisotropy

Magnetoelastic anisotropy describes the coupling between lattice strain/stress and magnetization. The magnitude of the effect is given by the magnetostriction \( \lambda \) which indicates the elongation of the crystalline lattice \( \Delta l(H) \) in an external magnetic field \( H \)

\[ \lambda = \frac{\Delta l(H)}{l(H = 0)} \] (2.8)

The inverse effect, known as inverse magnetostriction, describes how an induced lattice strain \( \sigma \) affects the magnetization. In an isotropic materials, inverse magnetostriction produces magnetoelastic anisotropy [44]

\[ e_{me} = -\frac{3}{2} \lambda_1 \sigma \left( \cos^2 \phi_M - \frac{1}{2} \right) \] (2.9)

where \( \phi_M \) is the angle between the directions of lattice strain and magnetization.

In cubic crystals, the magnetoelastic anisotropy energy can be written as [92]

\[ e_{me} = B_1 (c_1 a_1^2 + c_2 a_2^2 + c_3 a_3^2) + B_2 (c_4 a_2 a_3 + c_5 a_1 a_3 + c_6 a_1 a_2) + \ldots \] (2.10)

where \( B_i, a_i \) and \( c_i \) are the magnetoelastic coupling coefficients, cosines of the magnetization angles and strains with respect to cubic axis, respectively. The magnetoelastic coupling coefficients can be defined with magnetostriction constants \( \lambda_i \) and elastic stiffness constants \( c_i \)

\[ B_1 = -\frac{3}{2} \lambda_{111}(c_{11} - c_{21}) \] (2.11)

\[ B_2 = -3\lambda_{111}c_{44}. \] (2.12)

Here, the elastic stiffness is a tensor quantity which can be written using the
Voigt formalism for cubic crystals as
\[
c = \begin{pmatrix}
c_{11} & c_{12} & c_{12} \\
c_{12} & c_{11} & c_{11} \\
c_{12} & c_{12} & c_{11} \\
c_{44} & c_{44} & c_{44}
\end{pmatrix}
\] (2.13)

where \(c_{11}, c_{21}\) and \(c_{44}\) are Young’s modulus, transverse expansion and shear modulus, respectively. A lengthy expression, similar to Eq. 2.15, can be written for the effective field arising from this anisotropy [109].

### 2.5.4 Uniaxial anisotropy

In vectorial format, the energy density for uniaxial anisotropy up-to second order can be written as [109]
\[
e_{uni} = -K_u1 (\mathbf{u} \cdot \mathbf{m})^2 - K_u2 (\mathbf{u} \cdot \mathbf{m})^4
\] (2.14)

where \(K_1\) and \(K_2\) are the first and second order uniaxial anisotropy constants and \(\mathbf{u}\) is a directional vector. Uniaxial magnetic anisotropy can arise from magnetoelastic anisotropy via uniaxial strain, the crystalline structure or the shape of a patterned element. For example, cobalt with a close-packed-hexagonal structure exhibits a preferential easy-plane along the crystalline c-axis [85, 24]. The second order term is usually small in polycrystalline or amorphous materials but it can be important in ultrathin films [101]. By calculating the functional derivative, the effective field arising from uniaxial anisotropy is
\[
B_{anis} = \frac{2K_u1}{M_s} (\mathbf{u} \cdot \mathbf{m}) \mathbf{u} + \frac{4K_u2}{M_s} (\mathbf{u} \cdot \mathbf{m})^3 \mathbf{u}.
\] (2.15)

which introduces a Zeeman-like term to the total effective field.

### 2.5.5 Stoner-Wohlfarth model

Figure 2.1. Schematic illustration of the angles used in the Stoner-Wohlfarth model.
Hysteresis loops of a system with uniform magnetization and one preferential axis of magnetization (i.e. uniaxial magnetic anisotropy) are described by the Stoner-Wohlfarth model

\[ e = \frac{dE_{\text{tot}}}{dV} = K_u \sin^2(\varphi) - \mu_0 H M_s \cos(\varphi - \theta) \]  

(2.16)

where \( K_u \), \( \varphi \), \( \theta \), \( \mu_0 \) and \( M_s \) are the uniaxial anisotropy constant, direction of magnetization, direction of external magnetic field, vacuum permeability and saturation magnetization. The angles are defined with respect to the easy magnetization axis (see Fig. 2.1). This model can be used to extract the anisotropy strength from the hard-axis hysteresis loops (\( \theta = 90^\circ \)). The energy minima for this geometry are given by

\[ \frac{de}{d\varphi} = 2K_u \sin \varphi \cos \varphi + \mu_0 H M_s \sin(\varphi - \theta) = \cos \varphi(2K_u \sin \varphi + \mu_0 H M_s) = 0. \]  

(2.17)

The solution for the energy minima between the saturation fields (\( \frac{d^2e}{d\varphi^2} > 0 \)) is

\[ m_y = \cos(\pi/2 - \varphi) = \sin \varphi = \frac{\mu_0 M_s}{2K_u} H = kH. \]  

(2.18)

Thus, the magnetization component along the y-axis varies linearly with the strength of applied magnetic field. In experiments, the slope of the hard-axis hysteresis loop \( k \) can be used directly to extract the anisotropy constant

\[ K_u = \frac{\mu_0 M_s}{2k}. \]  

(2.19)

### 2.6 Magnetic domains

Exchange interactions in ferromagnetic materials prefer parallel alignment of neighbouring moments. However, even though the nearby moments are parallel at microscopic scale, the system can subdivide into regions with uniform magnetization but different orientation of magnetization between the areas (Fig. 2.2). These areas are known as magnetic domains. The formation of domains is induced by the competition between long-range magnetostatic interactions and short-range exchange interactions. The dipole-dipole interaction between the moments produces stray fields that are minimized by the formation of magnetic domains.

Magnetic domains are separated by magnetic domain walls with finite width. Inside the domain wall, the magnetization rotates gradually. The finite width of the wall is determined by a competition between exchange, magnetostatic and anisotropy energies \[ \ldots \]. Slow rotation of magnetization reduces the exchange energy at the expense of anisotropy energy.
Figure 2.2. Schematic illustration of an isotropic square ferromagnet at zero magnetic field. The system is subdivided into magnetic domains to minimize magnetostatic interaction at the expense of exchange energy. The circular field lines denote stray fields generated by the magnetization.
Ferromagnetism
3. Spin waves

At low temperatures, thermal energy is not sufficient to locally reverse magnetic moments (spins) due to the high cost in exchange energy (Eq. 2.3) \[24\]. However, the transverse components of magnetic moments can deviate from the equilibrium state and induce collective fluctuations of spins in the lattice (Fig. 3.1). These excited states are known as spin waves or magnons, the quanta of a spin wave. Spin waves can be divided into three wavelength scale regimes: exchange, dipole-exchange and dipolar spin waves. The first length-scale concentrates on spin-waves with wavelengths on the nanometer scale where exchange energy dominates and dipole-dipole interaction are negligible. By solving the Heisenberg Hamiltonian (Eq. 2.3), a quadratic dispersion relation can be derived for the angular frequency

\[
\omega = \frac{2Ja^2}{\hbar}k^2
\]

where \(s\), \(a\), \(\hbar\) and \(k\) are the spin, interatomic spacing, reduced Plank constant and wave vector, respectively. These highly energetic waves are usually probed using neutron \[74\] or inelastic x-ray scattering \[2\] which are valuable methods to estimate the exchange stiffness. At the other end of the wavelength spectrum, neighbouring magnetic moments are nearly parallel and the propagation of spin waves is mediated via long-range dipole-dipole interactions. Unlike exchange spin waves, the transport of dipolar spin waves is highly anisotropic, depending
Spin waves sensitively on the angle between the wave vector and the direction of magnetization [48]. Lastly, dipole-exchange spin waves are an intermediate lengthscale where both exchange and dipole-dipole interactions are relevant.

### 3.1 Landau-Lifshitz-Gilbert equation of motion

![Figure 3.2. Schematic illustration of the vectors in the Landau-Lifshitz-Gilbert equation.](image)

When a magnetic system is in an excited state, the torque extorted on the moments (Eq. 2.2) is nonzero. In this case, the excess energy is dissipated via a precessional motion around the equilibrium orientation. This dynamic motion was first described by Landau and Lifshitz using a classical variational procedure [66]. However, the derived equation of motion was insufficient to describe systems with strong dissipation [39]. Gilbert introduced a reformulated version of the Landau-Lifshitz equation by introducing a viscous-like damping term. In this form, the torque (Eq. 2.2) remains constant even when the damping factor is increased. In the Landau-Lifshitz-Gilbert (LLG) formalism, the magnetic torque $\mathbf{r}$ that is applied to a reduced magnetization $\mathbf{m}$ is

$$
\frac{dm}{dt} = \mathbf{r} = -\gamma \mathbf{m} \times \mathbf{B}_{\text{eff}} + \alpha \gamma \mathbf{m} \times \frac{dm}{dt}
$$

(3.2)

where $\gamma$, $\alpha$ and $B_{\text{eff}}$ are the gyromagnetic ratio, damping factor and effective magnetic field, respectively. The effective field includes contributions from magnetic anisotropy, exchange interaction, Zeeman interaction and magnetostatic interaction. The first term describes precessional motion around the effective field and the latter term introduces damping towards the effective field. The vectors in this equation are illustrated in Fig. 3.2. The main relaxation paths are spin-orbit coupling (magnon-phonon scattering or spin-lattice relaxation), magnetic impurities, magnon-magnon interaction and eddy currents. The LLG equation can be linearized in the absence of damping and solved to calculate the angular frequency of uniform precessional motion in thin films. This formula is known as the Kittel equation [56]

$$
\omega = \gamma \sqrt{B_{\text{eff}}(B_{\text{eff}} + M_s)}.
$$

(3.3)
To derive formulas for spin waves with finite wavelengths, the LLG equation needs to be combined with Maxwell equations and electromagnetic boundary conditions.

### 3.2 Dispersion relations

#### Figure 3.3.

a) Spin wave dispersion relations for a CoFeB film for various propagation directions with respect to the direction of in-plane magnetization. The system was modeled at 30 mT with zero magnetic anisotropy. b) Comparison between Damon-Esbach (Eq. 3.8) and backwards volume (Eq. 3.7) mode against the model by Kalinkos and Slavin (Eq. 3.9). At large wave vectors, the quadratic dispersion of exchange spin waves is evident.

The dispersion of spin waves is highly non-linear and anisotropic unlike electromagnetic radiation. The dispersion depends strongly on the angle between the wave vector $k$ and the direction of magnetization (Fig. 3.3a). Perpendicular or parallel configurations are typically used in experiments which are known as Damon-Esbach (DE), backwards volume spin wave (BVSW) and forward volume spin wave (FVSV) modes. The first two are in-plane configurations where the wave vector is perpendicular or parallel to the magnetization, respectively. The latter describes an out-of-plane magnetized film with spin wave propagating in-plane.

#### 3.2.1 Forward volume modes

The dispersion relation of spin waves in perpendicularly magnetized thin films is

\[
\omega^2 = \omega_H \left[ \omega_H + \omega_M \left( 1 - \frac{1 - e^{-kL}}{kL} \right) \right].
\]  

(3.4)

where

\[
\omega_M = \gamma \mu_0 M_s
\]  

(3.5)

\[
\omega_H = \gamma \mu_0 H_i
\]  

(3.6)
Spin waves

and $L$ is the film thickness. In this case, the magnetization is always perpendicular to the film leading to an isotropic in-plane dispersion relation. The isotropy is beneficial when spin waves are transmitted through translationally broken areas such as curves in spin majority gates [57].

3.2.2 Backwards volume mode

The dispersion relation of spin waves in in-plane magnetized films with wave vector parallel to the direction of magnetization is

$$\omega^2 = \omega_H(\omega_H + \omega_M(1 - e^{-kL}/kL)). \quad (3.7)$$

The name, backwards volume mode, stems from the sinusoidal mode profile across the film thickness, negative group velocity and positive phase velocity. Furthermore, BVSWs are reciprocal, i.e., the dispersion relation remains identical when the direction of the effective magnetic field reverses.

3.2.3 Damon-Eshbach mode

The dispersion relation of spin waves in in-plane magnetized films with wave vector perpendicular to the direction of magnetization is

$$\omega^2 = \omega_H(\omega_H + \omega_M) + \frac{\omega_M^2}{4}(1 - e^{-2kL}). \quad (3.8)$$

In contrast to forward and backward volume modes, the mode profile is concentrated on one of the film surfaces and decays exponentially into the bulk part of the film. If the field or propagation direction is reversed, the surface of propagation is inverted leading to non-reciprocity. This phenomenon is important in spin wave experiments because most experimental techniques probe only the top surface of a sample.

3.2.4 Dispersion relation for arbitrary propagation direction

The spin wave spectra discussed in the previous sections omitted exchange interactions between spins and considered only perpendicular or parallel orientations between the wave vector and the direction of magnetization. Kalinkos and Slavin derived a formula for exchange-dipolar spin waves with mixed boundary conditions for arbitrary propagation directions [57]. In their model, the pinning of spins at the two surfaces can be different and the strength can vary continuously between fully pinned and totally unpinned spins. The dispersion of such spin waves is

$$\omega^2 = (\omega_H + \alpha \omega_M k_n^2)(\omega_H + \alpha \omega_M k_n^2 + \omega_M F_{nn}) \quad (3.9)$$

where the dipole-dipole matrix element for the first order mode is

$$F_{00} = P_{nn} + \sin^2 \theta \left(1 - P_{nn}(1 + \cos^2 \phi) + \omega_M \frac{P_{nn}(1 - P_{nn} \sin^2 \phi)}{\omega_H + \alpha \omega_M k_n^2} \right). \quad (3.10)$$
For completely pinned surface spins, $P_{nn}$ can be written as

\[ P_{00} = 1 - \frac{1}{kL}(1 - e^{-kL}). \]  

(3.11)

The results for the arbitrary propagation directions reproduce accurately the dispersion of conventional Damon-Eshbach and BVSW modes at small wave vectors (Fig. 3.3b). However, the omission of exchange coupling in the expressions of DE and BVSW modes (Eq. 3.7 and Eq. 3.8) is evident at large wave vectors where their dispersion saturates. In contrast, the inclusion of exchange coupling in the model by Kalinkos and Slavin correctly introduces the quadratic nature of exchange spin waves at large wave vectors [48].

### 3.2.5 Spin wave group velocity

The group velocity of spin waves is defined as

\[ v_g = \frac{d\omega}{dk}. \]  

(3.12)

From the dispersion relations, it can be calculated that the group velocity (i.e. the slope of the spin wave dispersion curve) increases with saturation magnetization and film thickness. Thus, large magnetic moments and thick films are generally favorable for high spin wave group velocities.

### 3.3 Spin transfer torque

![Figure 3.4](image.png)

**Figure 3.4.** Schematic illustration of spin transfer torque effect. Spin polarized current from the left side extorts a magnetic torque on the magnetic domain wall.

Charge carriers can extort a small torque to the magnetization in spatially non-uniform spin textures and magnetic (tunnel) junctions. This phenomenon is known as spin transfer torque and it is based on transfer of angular momentum of the charge carrier to the magnetization (Fig. 3.4). Various applications of this torque have been proposed, such as current driven domain wall motion and magnetic memory cells where a strong current in a tunnel junction can induce a magnetic switching event. However, the applied torque is relatively small and highly localized, and thus, large current densities on the scale of $10^{10}$ A/m$^2$ are
Spin waves are generally required to create sufficient torque for magnetic domain wall motion or magnetic switching.

Spin transfer torque can be formally introduced to the Landau-Lifshitz-Gilbert equation by implementing additional terms arising from adiabatic and non-adiabatic torques. In Zhang and Li's formalism, when a current is applied to a slowly and adiabatically varying magnetization such as in a domain wall [89], the LLG equation is modified as [109, 120]

\[
\frac{dm}{dt} = -\gamma m \times H_{\text{eff}} + \alpha \gamma m \times \frac{dm}{dt} - \beta j \cdot \nabla m \quad (3.13)
\]

where \( j \), \( \beta \) and \( b_j \) are the current density, degree of non-adiabaticity and the coupling strength between the applied current and the magnetization, respectively. This approximation is valid when the spin diffusion length is shorter than the spin rotation of the non-uniform magnetic texture. The origin and strength of the non-adiabatic contribution is still under debate even though it plays an important role in STT induced domain wall motion and depinning [12]. In general, the non-adiabatic torque is highly sensitive to the texture of the domain wall. Small \( \beta \) values prefer domain wall propagation via nucleation of vortices and anti-vortices while larger values support driven motion without vortice nucleation [105].

### 3.4 Excitation of short wavelength spin waves

Initial experiments on ferromagnetic resonances were performed in a microwave cavity where the magnetic field of a standing electromagnetic field induces a finite torque to the magnetization [40, 54]. In this system, the oscillating magnetic field is uniform. To excite spin waves with a finite wavelength, a localized torque is required. A common excitation technique relies on Oersted fields created by a microwave structure such as a microstrip antenna or coplanar waveguide. In this case, the wavelength of the excited spin wave is roughly twice the width of the signal line. However, when the conductor width is reduced, Joule losses increase [86] and impedance matching between the antenna and microwave source becomes challenging. Impedance matching is required to minimize reflections of the electromagnetic wave from the microwave structures. Currently, the shortest excited wavelength from an antenna is 370 nm using a 125 nm wide antenna [23]. Various methods have been proposed for the excitation of spin waves with short wavelengths using large microwave antennas. One technique is based on a magnetic grating coupler where a periodic array of magnetic nanostructures is grown onto a magnetic film. In such a system, the reciprocal lattice vector of the array enhances the momentum of the emitted spin waves. For cobalt stripes on a yttrium iron garnet film, the excitation of spin waves with a wavelength below 100 nm has been demonstrated [71, 119]. Alternative techniques rely on dynamic fluctuations of a topological defects such as a vortex core [117] or magnetic domain wall (Publication IV). When a domain
Spin waves

wall is driven into resonance using an uniform ac-magnetic field [41, 113] or current [8], the DW radiates spin waves at the same or twice the frequency as the actuating signal. Spin waves can also be excited by coupling two forced magnetization precessions with different amplitudes using uniform microwave fields. Such excitation is realized in an anisotropy modulated BaTiO$_3$/CoFeB system (Publication V) and in exchange coupled YIG/CoFeB bilayers (Publication VIII). Dynamic demagnetization field near geometrical boundaries can also trigger excitation of propagating spin waves in semi-infinite films [27] or rectangles [81, 73] at an arbitrary excitation frequency.

3.5 Spin wave confinement

Eq. 3.9 describes the dispersion relation of an infinite film. However, physical boundaries are relevant in experiments where edges cause additional electromagnetic constraints [87]. For example, if a magnetic field is applied parallel to the edges of an elongated wire, the spins at the boundaries are partially pinned which leads to lateral quantization of spin waves. In case of fully pinned spins, the quantized wave follows a sinusoidal profile where the nodes exist at the geometrical boundaries [29]. On the other hand, quantization can be induced also in structures without physical boundaries if the confinement is induced by magnetic non-uniformities and/or variations of the effective magnetic field. For example, it has been shown that non-uniform magnetization textures can form a hundreds of nanometers wide magnonic waveguide, similar to an optical fibers, in zig-zag magnetization configuration of a wire [105] or inside a domain wall [114]. Confinement can also be induced by modulation of magnetic anisotropy in continuous films via ion irradiation [107] or via strain-coupling to ferroelastic domains of a ferroelectric crystal (see Publication V). Mathematically, the formation of a quantized mode can be described by a Bohr-Sommerfeld integral [46, 94]

$$\int_{x_0}^{x_1} k(x)dx = n\pi$$  \hspace{1cm} (3.14)

where $x_i$ are the turning points and $n$ is the order number. This quantization model is in good agreement with experiments on inhomogeneously magnetized wires [10].

3.6 Magnonic crystals

Magnetic resonances have been used in microwave engineering as signal filters [77] and signal sources [106] where the properties can be tuned by an external magnetic field. Conventional approaches exploit the resonance of yttrium iron garnet spheres that are coupled to a microwave circuit. However, macroscopic spheres are not easily implemented in complex and small spin wave devices.
Spin waves

Artificial metamaterials, so-called magnonic crystals, allow filtering at smaller length scales. In these materials, the magnetic properties are periodically modulated which induces band gaps. The lateral variation of magnetic properties can be induced by various techniques. For example, etched grooves \[20\] or holes \[83\] introduce additional electromagnetic boundary conditions. Alternatively, the saturation magnetization can be reduced via ion-implantation or a thermal profile \[112\]. Oersted fields \[18, 17\] or spin-orbit coupling in MgO/Co bilayers \[115\] has also been exploited. In these structures, a propagating spin wave interacts with the periodic lattice and the wave is strongly reflected if the Bragg condition

\[ n \lambda = 2d \sin \theta, \]

where \( n, \lambda, \theta \) and \( d \) are an integer, wavelength, angle between lattice and propagation direction, and lattice constant, is fulfilled. Bragg scattering induces band gaps in the dispersion relation that can be used in band pass filters. By tuning the periodicity and width of the disturbances, the width and strength of the forbidden gaps can be tuned \[21\]. In general, stronger modulation of magnetic parameters enhances the blocking strength at the cost of parasitic losses in the transmission band. For example, it has been shown that a thickness reduction of 1/10 in 5.5 \( \mu \)m thick YIG films is optimal to induce a strong rejection band and -3 dB insertion loss at the transmission band \[19\]. If a single filter frequency is required, higher order rejection bands can be suppressed by varying the magnetic parameters harmonically.

3.7 Low damping materials for magnonics

Materials with small Gilbert damping factor and long propagation lengths are desirable for spin wave devices \[22\]. Currently, one of the best candidates is single crystal yttrium iron garnet (YIG) which exhibits the lowest damping factor of about \( 10^{-5} \) in micrometer thick films \[68\]. However, the properties of YIG are extremely sensitive to the crystal quality due to its complex crystalline structure with a large unit cell size of 12.4 Å that contains 80 atoms \[14\]. High quality crystals have been grown by liquid phase epitaxy on (111) oriented gadolinium garnet substrates. This technique is most suitable for micrometer thick films. Recent advancements in epitaxial growth have also enabled the deposition of sub-micrometer YIG films with near bulk-like quality \[32\]. Further reduction of the film thickness to sub-100 nm is important for spin wave devices because patterning thicker films with nanometer scale feature size is challenging. Recently, the growth of nanometer thick YIG films by pulsed laser deposition \[102\] and magnetron sputtering \[72\] has attracted a lot of attention. For example, a damping factor of \( 8.5 \times 10^{-5} \) has been reported in a 22 nm thick YIG film \[102, 13\]. The drawback of YIG is its high deposition and/or annealing temperatures of >500°C \[60\] which is not compatible with conventional comple-
mentary metal-oxide semiconductor fabrication processes [96]. Furthermore, the saturation magnetization of YIG is rather small which limits its operation frequency and propagation velocity (Eq. 3.12). Other good materials for magnonics are Permalloy and CoFeB alloys which exhibit damping factors of the order of $10^{-3}$. Because these metallic films do not require epitaxial growth, they are easily fabricated by electron beam evaporation or magnetron sputtering at room temperature. Their saturation magnetization is larger than YIG and their magnetic anisotropy tends to be weak.
Spin waves
4. Multiferroic heterostructures

Multiferroics are materials which exhibit multiple ferroic properties, ferromagnetism, ferroelectricity and/or ferroelasticity, simultaneously [90]. Recently, the term has been expanded to incorporate also antiferromagnetism and ferrimagnetism. The interest in multiferroics has risen in the past decades due to potential applications in actuators, filters and sensors [43]. Especially the coupling between ferroelectricity and ferromagnetism, known as magnetoelectric multiferroics, is intriguing for applications [110]. Such systems could be exploited in magnetic memories where the magnetic state is written electrically and read-out magnetically [7] or both non-volatile phases could be used to store information in a four-state memory [38].

Multiferroic materials can be divided into two categories: single-phase and heterostructures [34]. Single-phase multiferroic materials exhibit multiple ferroic orders simultaneously. In most single-phase multiferroics, at least one of the ordering temperatures (magnetic or electric) is low and coupling between the two ferroic phases tends to be weak [43]. One commonly studied material is BiFeO$_3$ which exhibits ferromagnetic coupling between the Fe atoms along the pseudo-cubic (111) direction and antiferromagnetic coupling between adjacent planes. However, the antiferromagnetic spin spiral axis rotates through the bulk crystal over tens of nanometers leading to zero macroscopic moment [99]. An alternative approach relies on magnetoelectric coupling in multiferroic heterostructures where each ferroic phase can be engineered for high-temperature operation and strong magnetoelectric coupling can be obtained via interactions at their interface [34, 75]. The coupling can be based, for example, on exchange coupling at the interface between an antiferromagnet/ferroelectric and a ferromagnet or strain transfer from ferroelastic domains in an ferroelectric (Publication I-VII). Both coupling mechanisms can produce large modulations of magnetic anisotropy in the ferromagnetic film, opening up ways to control magnetism via the application of electric fields [90, 63, 15].
4.1 Domain pattern transfer in BaTiO$_3$-based multiferroic heterostructures

Spatially modulated anisotropy patterns can be achieved by top-down or bottom-up processes. For example, a top-down lithographic technique relies on patterning of an exchange-biased IrMn/Py bilayer using ion irradiation [107] or thermal scanning probe lithography [1]. In the first case, a spatially uniform unidirectional exchanged bias is initialized during sample growth followed by a mask patterning process. The unmasked area of the surface is bombarded by He ions under application of an external magnetic field which locally rotates the direction of exchange bias leading to an anisotropy modulated pattern. The latter technique is based on applying a magnetic field opposite to the initial exchange bias direction while using a atomic force microscopy tip to locally heat the exchange biased structure above the Néel temperature. This technique allows nanoscale control of the direction of unidirectional anisotropy axis. Similar anisotropy modulated landscapes are also achieved in a bottom-up process using ferromagnetic-ferroelectric multiferroic heterostructures (Publications I-VII) [64]. In this case, the anisotropy pattern is induced by strain transfer from a ferroelectric substrate with regular ferroelastic stripe domains [34, 61, 36].

In this thesis, single-crystal ferroelectric BaTiO$_3$ substrates are used to fabricate strain-coupled multiferroic heterostructures with regular modulations of magnetic anisotropy. BaTiO$_3$ exhibits a tetragonal crystalline structure at room temperature and the spontaneous ferroelectric polarization is oriented along the crystalline c-axis [49]. If an electric field is applied perpendicular to the polarization, an abrupt 90° rotation of the polarization axis coincides with a 90° rotation of the lattice tetragonality [63]. This characteristic feature of BaTiO$_3$ allows switching between two non-volatile strain states and, consequently, control of magnetic anisotropy via inverse magnetostriction.

BaTiO$_3$ (001) substrates exhibits three highly symmetric ferroelectric domain configurations. In 180° ferroelectric domains, the polarization is oriented in out-of-plane direction. In the a-c domain configuration, the direction of ferroelectric polarization alternates between in-plane (a-domains) and out-of-plane (c-domains). In the in-plane polarized stripe domains, the polarization is oriented perpendicular to the boundary. In the so-called a1-a2 configuration, the ferroelectric polarization, and thus, the tetragonal c-axis, rotates by 90° at the boundaries between elongated stripe domains. The polarization inside the
Multiferroic heterostructures

domains makes an angle of 45° with respect to the domain boundaries.

Strain coupling between ferroelastic domains of a BaTiO3 substrate and a ferromagnetic film has been achieved for polycrystalline CoFe films [64, 61, 66], epitaxial Fe films [114, 62], La1-xSrxF3 [69, 33] and partly amorphous films (e.g. CoFeB [5, 10, 45] and Publications V and VII). In this thesis, two experimental systems are used to study domain-wall and spin-wave dynamics in multiferroic heterostructures. In publication IV, a 20 nm thick iron film was grown on top of a BaTiO3 substrate with an a-c ferroelastic domain configuration using molecular beam epitaxy at 300°C. The strain coupling between the rectangular unit cell of in-plane a-domain of the BaTiO3 substrate induce uniaxial anisotropy in the Fe film. The magnetic easy axis is parallel to the stripe domain boundary due to negative magnetostriction of Fe. On top of the square c-domains, the anisotropy is fourfold symmetric [62]. Ferroelectric polarization at the boundary rotates within few unit cells [121] and, consequently, the magnetic anisotropy boundary is abrupt and magnetic domain walls are strongly pinned to the boundaries.

In Publication II, we demonstrated that strong pinning of magnetic domain walls onto ferroelectric boundaries could be utilize to reversibly drive the magnetic domain walls by an electric field without any assistance of a magnetic field or electric current. In the experiments, the multiferroic system is first initialized by a magnetic field and then the field is turned off. Next, electric field pulses are applied across the BaTiO3 substrate using a back electrode and the Fe film as top contact. Depending on the polarity of the field, the ferroelastic c-domain grows/shrinks at the expense of the a-domain. Lateral motion of the ferroelectric domain walls also moves the magnetic anisotropy boundaries in the Fe film. As a result, the pinned magnetic domain walls are purely driven by the electric field.

Publications III, V and VII, the experimental multiferroic heterostructures consist of a CoFeB film (40% Co, 40% Fe, 20% B) on top of a BaTiO3 substrate with an a1-a2 domain configuration. To realize full domain pattern transfer, the substrate was heated to 175° or 300°C during CoFeB film growth. At these temperatures, the BaTiO3 substrate is in paraelectric and its crystalline structure is cubic [78]. After sputtering of a 50 nm thick CoFeB layer, the sample was cooled down to room temperature. Upon cooling through the ferroelectric transition temperature at about 120°C, the substrate retains its original tetragonal lattice structure. As a result, an a1-a2 domain pattern forms. At room temperature, the lattice tetragonality has grown to 1.1% [49, 78]. Part of this strain (typically 10-25%) is transferred to the CoFeB film [61]. Via inverse magnetostriction, it induces a modulation of magnetic anisotropy, whereby the uniaxial magnetic easy axes are oriented parallel to the in-plane direction of ferroelectric polarization. CoFeB was chosen in Publications III, V and VII because it combines large magnetostriction [67] with relatively low magnetic damping [25, 70]. Because of this, the magnetic anisotropy of the CoFeB film is dominated by magnetoelastic contributions and domain formation within stripe domains is avoided.
Multiferroic heterostructures
5. Methods

The samples for my thesis project were grown using molecular beam epitaxy, magnetron sputtering and pulsed laser deposition. The samples were patterned using photo or electron beam lithography depending on the required feature size. The magnetostatic properties were characterized using magneto-optical Kerr microscopy and the magnetodynamic properties were measured using Brillouin light scattering and all-electrical spin wave spectroscopy. The results from the experiments were compared to micromagnetic simulations to interpret the underlying physics.

5.1 Sample preparation

5.1.1 DC magnetron sputtering

Sputtering is a physical vapor deposition technique in which a target material is bombarded with energetic ions or atoms. The collision induces ballistic ejection of clusters from the target material due to momentum transfer. The bombardment is achieved by injecting a process gas, usually argon, to the vacuum chamber and applying a negative voltage to the target. When the electric field ionizes a gas molecule, it is accelerated towards the target where the collision releases near-stoichiometric clusters of the target material. To improve the plasma stability and enhance the ionization probability via collisions between gas particles, permanent magnets are added underneath the target \cite{50}. The magnetic field gradient partially confines electrons into a volume where a high density of charged particles with long free-path length can collide with the neutral process gas. This technique allows significantly lower growth pressure and higher deposition rates. However, the localization of the plasma leads to non-uniform utilization of the target material. Sputtering can also be used to grow oxides and nitrides from metallic targets \cite{82}. When a reactive gas is injected to the chamber, the bombardment causes a chemical reaction with the surface of the target. This reactive sputtering technique allows for reasonably fast deposition.
rates of oxides.

In this thesis, the BaTiO$_3$/CoFeB samples were prepared using magnetron sputtering. The samples were mounted on a 2” silicon wafer using silver paste. Prior to deposition, the samples were heated to 175$^\circ$ or 300$^\circ$C using high intensity halogen lamps behind the wafer. The temperature was increased at a rate of 10$^\circ$C per minute and stabilized at the final temperature for roughly ten minutes. The 50 nm thick CoFeB film was grown at 6.6 mbar Argon pressure and 50 W power leading to 0.12 nm/s deposition rate. After the sample was cooled down to room temperature, a 5 nm Au capping layer was grown to prevent oxidation of the magnetic film (Publications III and V). In Publication VII, the sample was capped with a 3 nm thick tantalum layer followed by reactive sputtering of 28 nm TaO$_x$ insulating layer to prevent shorting of the excitation antennas on top of the sample.

### 5.1.2 Pulsed laser deposition

Pulsed laser deposition is a versatile technique that is suitable for a wide variety of materials. This technique relies on the interaction between a temporally short and high energy laser pulse and the target material. When the pulse is absorbed by the target, the surface layer rapidly melts and vaporizes into a neutral and/or charged gas plume that follows closely the composition of the material. However, light elements, such as oxygen, exhibit a different angular ejection distribution than heavier elements. To fully control the composition, films can be grown in, e.g., an oxygen atmosphere or by using targets with slightly different stoichiometry than the desired in final film.

In publication VIII, 80 nm and 295 nm thick YIG films were grown on single-crystal GGG(111) substrates using pulsed laser deposition. The substrates were degassed at 550$^\circ$C for 15 minutes. After 0.13 mbar oxygen pressure was added to the deposition chamber, the substrate temperature was increased to 800$^\circ$C at a 5$^\circ$C per minute rate before deposition. The stoichiometric YIG target was exposed to an excimer laser with a pulse repetition rate of 2 Hz and a laser fluence of 1.8 J/cm$^2$. The films were post-deposition annealed at 730$^\circ$C in 13 mbar oxygen environment for ten minutes. Hereafter, the samples were cooled back to room temperature at a -3$^\circ$C per minute rate.

### 5.1.3 Lithography

Photolithography and electron beam lithography are microfabrication techniques to pattern films. The former is generally used to pattern structures at large production scales. In contrast, the latter is a maskless and low throughput technique that is generally used in prototyping and mask fabrication with nanometer resolution. In both techniques, a light or electron-beam sensitive resist is spin coated on a flat substrate at high rotation velocity followed by a baking process to evaporate excess solvent. The exposed areas of the resist are
Methods

Coat with resist Exposure Development

Sputtering Lift-off

PMMA MAA

Substrate Substrate Substrate

Ta/Au

Substrate Substrate

Figure 5.1. Illustration of a PMMA/MAA double-resist lift-off lithography process.

degraded or strengthened depending on the polarity of the resist. By masking the exposure or scanning a high intensity beam over the resist, the exposed area becomes selectively soluble or insoluble in a developing agent. After development, a patterned resist mask is formed on the sample surface. The mask can be used to selectively remove material via ion milling or chemical etching. Alternatively, material can be deposited on top of the sacrificial layer followed by a lift-off step where the resist, and the material on top of it, is removed using a solvent (Fig. 5.1). In some cases, the deposition material covers also the sidewalls of the resist, which can complicate the lift-off process. Furthermore, the sidewall material may remain contacted to the pattern and induce hanging ears that adhere to the surface. This effect can cause shorts in high density patterns. To overcome these issues, a double layer resist can be used in which the bottom layer is soluble to the developer. In this case, an undercut is formed during development and the deposited material does not adhere onto the sidewalls of the resist layer.

5.2 Magneto-optical Kerr microscopy

Spatially resolved measurements of magnetization can be performed using various techniques. For example, nanometer scale details can be resolved using scanning electron microscopy with polarization analysis [93] or x-ray magnetic circular dichroism [35]. However, these techniques are expensive and require large facilities. An inexpensive alternative is magneto-optical Kerr microscopy where the magnetization is probed using light. In this technique, linearly polarized light is reflected from the sample surface. During the reflection, the polarization state or light intensity is change depending on the orientation of
magnetization (Kerr effect) [88]. The rotation of polarization can be measured using a polarizer-analyzer configuration at near extinction angle. The resulting light is captured by a sensitive CCD camera to gain spatial resolution. The Kerr effect exist in three configurations: longitudinal, polar and transverse. In the first two effects, the light undergoes a change in polarization that is proportional to the angle between the optical axis and magnetization. If the change in ellipticity is more prominent than polarization rotation, a rotatable quarter wave plate can be introduced prior to the analyzer to improve sensitivity of the measurement. The transverse Kerr effect induces magnetization dependent intensity changes and it is generally weaker than polar or longitudinal Kerr effects. In addition to the Kerr effect, the polarization state is also influenced by a Faraday effect in the microscope objective. This effect produces a rotation of polarization that is linearly proportional to the applied magnetic field. This additional effect can be removed from the measurement by fitting a slope to the signal beyond saturation and subtracting it from the measurement data.

The magnetostatic properties presented in this thesis were measured by an optical reflection microscope from Zeiss that was modified for magnetic imaging by Evico Magnetics. In this system, high intensity light from a LED source is directed to a high NA-factor objective. To control the incidence angle, an adjustable aperture is added to the beam path. A finite incidence angle is selected by setting the aperture such that only one side of the objective is used.

![Optical image](image1)

**Figure 5.2.** Optical image and magnetic contrast image after background subtraction of a BaTiO$_3$/50 nm CoFeB sample.

Captured images of the magneto-optical Kerr microscope contains both magnetic information and optical details of the sample surface. In most cases, effects caused by surface roughness and defects dominate the optical contrast and magnetic information is not directly visible. To resolve magnetic domains, the optical background needs to be subtracted. Background subtraction is achieved by the acquisition and averaging of multiple images in an oscillating (∼10 Hz) triangular magnetic field. Because the field sweep is much faster than the averaging time, the magnetization is zero on average and the resulting image contains only optical information of the sample surface. By subtracting this background
image from the measured data, the magnetic contrast is significantly improved. The background subtraction process is illustrated in Fig 5.2.

In the microscope program, a pre-selected region can be chosen and the measured intensity is averaged. This allows measurement of localized hysteresis loops. The local measurement technique was used in Publications I-III, V and VII to extract the magnetic anisotropy strength, location and width distribution of magnetic stripe domains, and/or to confirm successful imprinting of the anisotropy modulated pattern.

5.3 Brillouin light scattering

Brillouin light scattering (BLS) is one of the most sensitive techniques to measure spin waves [30]. Unlike all-electrical spin wave spectroscopy where the excited wave-vector distribution is fixed by the antenna [118], the measured wave vector can be chosen and continuously tuned in BLS. Thus, full dispersion relation can be extracted. BLS relies on the inelastic scattering of photons from magnons. During the scattering process, a magnon can be created or annihilated while the energy and momentum $k$ of the incident photon and magnon are conserved. In the case of a creation process, the scattered photon transfers momentum and energy to the magnon [30]

$$\hbar k_i = \hbar (k_r + k_m) \quad (5.1)$$

$$\hbar \omega_i = \hbar (\omega_r + \omega_m) \quad (5.2)$$

where the subscripts, $i$ and $r$, correspond to incident and reflected photons, respectively. In thin films, where the wavelength of light is significantly larger than the thickness of the film, only in-plane components of the photon momentum must be considered. In backscattering geometry, an incident laser beam

![Figure 5.3. a) Typical thermal Brillouin light scattering spectrum of SrTiO3/50-CoFeB measured at 100 mT and an incident angle of 50°. b) Dispersion relation of SiO2/20 nm CoFeB measured at 100 mT in-plane magnetic field and in Damon-Eshbach configuration. In this experiment, the incident angle was tuned between 10 and 65°. In the spectra, both perpendicular standing spin wave (PSSW) and the Damon-Eshbach dispersions are visible.](image)

grazes the sample surface at a finite incident angle $\theta$ and the backscattered photons are detected. The transferred in-plane momentum is

$$k_m = 2k_i = \frac{4\pi}{\lambda} \cos \frac{\theta}{2}$$  \hspace{1cm} (5.3)$$

where $\lambda$ is the wavelength of the photon. Because the transferred momentum depends on the incident angle, the wave vector of the spin waves is well defined. By measuring the energy loss of the photon

$$\hbar\omega_m = \hbar(\omega_i - \omega_r)$$  \hspace{1cm} (5.4)$$

$$f_m = f_i - f_r$$  \hspace{1cm} (5.5)$$

the magnon frequency $f_m$ is obtained. The change in frequency shift of the photon is measured using a tunable Fabry-Perot interferometer and a photodetector. Usually, two interferometers with different central frequencies and multiple passes are used to improve the sensitivity and to increase the measured frequency range. Typical BLS measurements are shown in Fig. 5.3.

BLS can also be used to obtain spatially resolved information. In this case, the laser is focused on the sample using an objective lens and the reflected light is collected. However, the focused light forms a large cone angle and, thus, wave vector resolution is lost. The wave vector can be retrieved by examining the phase of propagating spin waves that are excited by a microwave antenna. In this phase-resolved micro-focused BLS technique, the incident light is modulated with an electro optical modulator at the excitation frequency [97]. Interference of the reflected beam and a reference beam at the detector produces an intensity modulation that depends on the spin-wave phase.

5.4 All-electrical broadband spin wave spectroscopy

The experimental setup used to characterize magnetodynamics in this thesis utilizes ground-signal-ground microwave probes to connect coplanar waveguides to a vector network analyzer (Fig. 5.4). The waveguide is placed at the center of a quadropole magnet that can apply magnetic fields up to 270 mT in the film plane. The field strength of the magnet is continuously measured by two hall sensors underneath the sample stage and controlled using a PID loop to maintain constant magnetic field during longer measurements. The excitation antenna can be implemented into the experiment using two different techniques: either by patterning the structures on top of the sample or by placing the sample face-down on top of a prepatterned antenna. The latter flip-chip technique allows quick and simple characterization of samples with uniform magnetic properties across the sample. However, when measuring patterned samples, such as antidot lattices, it is often more convenient to directly pattern the antennas on top of the sample. Usually the antennas are optimized to an impedance of $Z_0 = 50 \Omega$ to minimize microwave losses and to allow easy implementation.
Methods

Microwave probe positioner
Quadropole magnet
Sample holder
Microscope
Vector network analyzer

Figure 5.4. All-electrical spin wave spectroscopy setup. The optical microscope and microwave positioners are used to connect probe tips to the microwave structures on the sample. The magnetic field strength is actively measured by two Hall sensors mounted below the sample stage. The field strength is actively controlled by a programmatic PID feedback loop.

of commercial signal sources that are optimized for 50 Ω impedance above gigahertz frequencies. The design of the coplanar waveguide also determines the excitation strength and wave vector microwave field. These parameters can be estimated using high-frequency electromagnetic simulations as shown in Fig 5.5. Furthermore, multiple waveguides with different signal line widths and known excitation spectra can be used to calculate the group velocity [10]. To extract the group velocity, Eq. 3.12 can be approximated as

$$v_g = \frac{2\pi df}{dk} \approx \frac{2\pi \Delta f}{\Delta k}$$

(5.6)

where $\Delta f$ and $\Delta k$ are the difference in resonance frequency and the main wave vector excitation microwave magnetic field of the waveguides.

In broadband spin wave spectroscopy, a high frequency rf-current in the signal line creates an oscillatory $h_{rf}$ Oersted field. This field exerts a local magnetic torque $\tau = h_{rf} \times M$ on magnetization $M$ in a nearby magnetic film which induces a frequency dependant response. At resonance, the magnetic field from the precessing moments induces a voltage in the signal line that is proportional to the complex susceptibility $\chi$. In all-electrical spin wave spectroscopy, the absorption spectrum is measured using a vector network analyzer that measures the transmission of a high frequency signal through the waveguide. However, the absorption is generally weak and buried underneath responses from properties unrelated to spin wave excitations, such as reflections and losses from the microwave circuitry (Fig 5.6a). The signal quality can be improve by calculating
the effective permeability of the sample \([47]\)

\[
U = \frac{i \ln[S_{21-m}(f)/S_{21-ref}]}{\ln[S_{21-ref}(f)]}
\]  

(5.7)

where \(S_{21-m}\) is the scattering parameter of the actual measurement and \(S_{21-ref}\) is the scattering parameter of a reference measurement. Spin wave excitations in the reference measurement are located outside the region of interest, which is usually realized by the application of a large external magnetic field. An example of an broadband spin-wave spectroscopy measurement in flip-chip geometry is shown in Fig. 5.6b.

The advantage of all-electrical broadband spin wave spectroscopy is its high sensitivity, frequency resolution (~1 Hz) and high measurement speed. However, the spin waves are probed over a macroscopic area and, thus, no spatially resolved information can be obtained. In order to gain insight into the spatial distribution of spin wave modes, the resonance frequencies are usually compared to micromagnetic simulations or theoretical models.

### 5.5 Micromagnetic modeling

Micromagnetic simulations are a useful tool to investigate magnetodynamic properties of complex magnetic systems. For small features in nanometer scale,
Methods

Figure 5.6. a) $S_{21}$ scattering parameter of the actual measurement (black) and reference measurement (red). In the experiment, SrTiO$_3$/50 nm CoFeB sample was measured using flip-chip technique at 30 mT magnetic field. The bias field was applied along the coplanar waveguide. b) Absorption spin-wave spectrum after background subtraction. The ferromagnetic resonance frequency is indicated with FMR.

Atomistic simulations can be used where each magnetic moment is discretely modeled. However, when the system size is sufficiently large, it becomes computationally challenging. In macroscopic size simulations, a continuum model is applied to the system where Landau-Lifshitz-Gilbert equation holds continuously. A GPU accelerated micromagnetic software, MuMax3 [109], solves the dynamic properties by discretizing the simulated area into finite cuboids. Each voxel is assumed to behave as a macrospin and magnetic interactions are considered as finite differences. This technique is suitable for systems with rectangular symmetry as rounded corners induce a stair-case effect which causes additional shape anisotropy. This effect can be avoided by smoothing the transition near rounded corners or by utilizing software packages which solves the LLG equation using tetrahedra as the finite element at additional computation cost. In the former case, the strength of the saturation magnetization is smeared near the edges to smoothen the transition.

5.5.1 Dynamic simulations

In thin ferromagnetic films, the magnetization in the ground state lies mainly in the plane of the film due to the demagnetization field. During excitation, the magnetic moments precess in an elliptical cone with a small magnetization component in the out-of-plane direction. Dynamic simulation spectra can be extracted by the spatially averaging the $z$-component of magnetization ($\langle m_z \rangle$). In dynamic simulations, the system relaxes to its ground state by minimizing the magnetic torque on the simulation mesh. After relaxation, the system is excited using small magnetic field pulses or continuous ac-excitation which mimics all-electrical spin wave spectroscopy. Location of resonance frequencies using continuous excitation requires very large computation times. In many cases it is advantageous to use pulse excitation for finding resonance frequencies because a broad range of frequencies up to a cut-off are excited simultaneously. In publication V, a short Gaussian field pulse was used to excite a wide fre-
quency range and the magnetization dynamics was simulated for 10 ns until the precessional motion was damped below numerical accuracy. The benefit of a Gaussian field pulse is good localization but its excitation strength decays as a function of frequency. Uniform excitation strength are achieved by sinc-type pulses. However, the magnetic field strength of these pulses decays as $1/t$ which increases the required simulation time. Results from both excitation methods are the identical but the relative amplitude of resonances in the spin wave spectrum differ.

From dynamic micromagnetic simulations, frequency spectra can be calculated by Fourier transforming the average $m_z$ component of the simulation mesh

$$A(f, B) = |\text{fft}(\langle m_z(t, B) \rangle)|. \quad (5.8)$$

The result represents the average spin wave amplitude as a function of frequency. During this process, spatial information magnetization precession is lost and even spin-wave modes are not resolved due to averaging. This issue can be avoided by saving the $m_z$ component of all cells at the expense of storage space usage, as discussed in the next chapter.

5.5.2 Spatially resolved simulations

Fourier imaging analysis allows visualization of localized and propagating spin waves [11, 105]. Instead of operating on the average magnetization, the transformation is performed separately for each simulation mesh cell

$$A(x, y, f) = \text{fft}(m_z(x, y, t)). \quad (5.9)$$

The localized spin wave modes can be visualized by presenting the magnitude of $A$ as a function of frequency and location in a contour plot. However, the phase information of the spin waves is lost and the plot illustrates only the intensity of precessional motion. As a consequence, the wavelength of propagating spin waves cannot be extracted. An alternative technique exploits continuous ac magnetic field excitation (Publications IV-VIII). Here, the resonance frequencies of the magnetic system are first extracted using pulsed excitation. After the system has returned to the relaxed ground state, a sinusoidal ac-magnetic field is introduced at the resonance frequency of the system. The time evolution of the continuously excited system can be studied spatially by saving the magnetization configuration at constant time steps. Advantages of this technique include the ability to gain insights into the spin wave formation processes, superior frequency resolution and sensitivity to the phase relationship between the excitation and magnetization precessions.
6. Results

This chapter discusses the excitation and propagation of spin waves in strain-coupled multiferroic heterostructures and CoFeB/YIG bilayers. The systems are investigated using all-electrical broadband spin wave spectroscopy and Brillouin light scattering on experimental samples in combination with detailed micromagnetic simulations.

Publications IV reports on spin wave emission triggered by forced oscillations of a pinned magnetic domain wall in an anisotropy modulated BaTiO$_3$/CoFe structure. The motion is driven by an ac spin-polarized electric current that pumps energy into the DW via a spin transfer torque effect. The system supports efficient excitation of short-wavelength spin waves up to high frequencies in magnetic nanowires where the lateral confinement produces an abundance of DW resonance modes. In publication VI, the static and dynamic properties of the pinned domain walls are theoretically modeled. The analytical framework predicts the domain wall width, mass and resonance frequency as function of anisotropy strength and external magnetic field. Publication V discusses confinement and emission of spin waves in a BaTiO$_3$/CoFeB multiferroic heterostructure at magnetic saturation, i.e., without magnetic domain walls. In this case, abrupt rotations of magnetic anisotropy induce a shift of the effective magnetic field at the anisotropy boundaries. The magnitude of this shift depends on the orientation of magnetization. Due to the disparity of the effective magnetic field in neighboring domains and, hence, a difference in the spin wave dispersion relations, a uniform microwave magnetic field excites higher order spin waves inside the anisotropy domains of the continuous ferromagnetic film. Furthermore, if a confined mode is generated inside domains with higher effective field, spin waves are emitted from the boundaries into the neighbouring stripe domains. The wavelengths of the propagating spin waves is significantly smaller than that of the confined mode and it can be continuously tuned by rotation of the magnetic bias field. A uniform microwave magnetic field also excites propagating spin waves from a single anisotropy boundary. This effect, which is caused by dynamic exchange coupling between different forced magnetization oscillations in neighboring domains, enables excitation of short-wavelength exchange-dominated spin waves. Finally, the multiferroic BaTiO$_3$/CoFeB sys-
tem is exploited as a reprogrammable spin wave filter in zero magnetic field. In this device concept, the pinned domain walls are switched between narrow head-to-tail and broad head-to-head or tail-to-tail configurations. The latter configurations are fully transparent to spin waves while the transmission of spin waves is significantly suppressed in narrow domain walls.

Publication VIII reports on the formation of perpendicular standing spin waves in YIG/CoFeB bilayers. Due to the large difference in saturation magnetization, the forced magnetization precessions inside the two layers are dissimilar. As a result, a dynamic exchange torque is generated at the interface and this triggers the launching of spin waves into the YIG film at the excitation frequency. If the wavelength of the emitted spin waves matches a perpendicular confinement condition, constructive interference produces a perpendicular standing spin wave mode.

6.1 Spin wave emission from an oscillating pinned magnetic domain walls

![Figure 6.1. Schematic illustration of the simulation are in the BaTiO$_3$/CoFe multiferroic heterostructure. The arrows illustrate the orientation of magnetization in the two 3.2 μm wide stripe domains.](image)

Publication IV concentrates on spin wave emission from a firmly pinned domain wall that is driven into oscillatory motion by a microwave spin-polarized current. As a model system, a multiferroic heterostructure consisting of a 15-nm-thick CoFeB film onto a BaTiO$_3$ substrate is considered [64, 61]. Experiments have shown that strain coupling between a1-a2 domains in BaTiO$_3$ and CoFe induces a regular modulation of magnetic anisotropy and strong pinning of magnetic domain walls onto the anisotropy boundaries. Because of pinning, the magnetic domains walls can be excited by an ac spin-polarized current via the spin transfer torque effect. Such oscillatory motion is expected to trigger the emission of spin waves [7, 8]. To test the prospect of localized spin-wave emission from a pinned magnetic domain wall, micromagnetic simulations were performed in MuMax3. In the simulations, domain pattern transfer from the BaTiO$_3$ substrate to the CoFe film was implemented by abrupt rotation of the uniaxial anisotropy axis at the ferroelectric domain boundaries. The simulation area of 6.4×3.2 μm$^2$ consisted of two anisotropy domains with an abrupt boundary at the center (Fig. 6.1). Periodic boundary conditions were applied in the film plane to mimic infinitely long stripes with widths of 3.2 μm. A saturation
magnetization of $1.7 \times 10^6$ A/m and exchange constant of $2.1 \times 10^{-11}$ J/m were used, in agreement with experiments [64, 61]. All simulations were performed in zero magnetic field after initialization of a 90° head-to-tail Néel domain wall.

To locate the domain wall resonance frequency $f_{res}$, a 1 ns long rectangular $5 \times 10^{12}$ A/m² pulse was first applied perpendicular to the domain wall. The magnetic response was simulated for 0.1 μs and the dynamic properties were investigated by averaging the magnetization over a 400 nm wide region around the domain wall followed by Fourier transformation. Simulation results for pulsed currents revealed that the magnetization inside the domain wall processes around the axis perpendicular to the domain wall, which corresponds to back-and-forth motion.

The potential stiffness $\kappa$ and domain wall mass $m$ were investigated as function of anisotropy strength $K_u$ by driving the oscillatory motion at the resonance frequency with an ac-current with gradually increasing current amplitude. The energy of the system was calculated each time the domain wall reached its maximum displacement. At this time, the velocity of the domain walls is zero and dynamic effects do not contribute to the total energy of the system. The potential stiffness $\kappa$ was extracted from

$$E = \frac{1}{2 \kappa x_c}$$

(6.1)

where $x_c$ is the domain wall displacement. The domain wall mass was calculated from $m = \kappa/2\pi f_{res}$. The simulation results were compared to an one dimensional theoretical model for a 90° domain wall. The model predicts the following exponential dependencies as a function of anisotropy strength $K_u$

$$f_{res} \sim K_u^{1/2}$$

(6.2)

$$\kappa \sim K_u^{3/2}$$

(6.3)

$$m \sim K_u^{1/2}$$

(6.4)

$$\Delta \sim K_u^{-1/2}.$$  

(6.5)

The simulation results were found to be in good agreement with the simulations. A more comprehensive analytical model of current-induced oscillations of a pinned 90°head-to-tail Néel domain wall is presented in Publication VI. In this paper, the model is compared to numerical simulations based on a 1D Heisenberg model and micromagnetic simulations.

Tunable spin wave emission from a the pinned magnetic domain wall was investigated by driving the domain wall by an ac spin polarized current with a density of $10^{12}$ A/m². The uniaxial magnetic anisotropy strength was set to $5 \times 10^4$ J/m³. An example is shown in Fig. 6.2a which depicts the profile of the out-of-plane magnetization component ($m_z$) after the onset of a 9 GHz current. The spin wave nodes that are radiated from the domain boundary are denoted with a solid black line after each excitation period and the slope corresponds to the group velocity of the emitted waves. The dotted envelope indicate the
Results

Figure 6.2. a) Simulated dispersion of emitted spin waves for $K_u = 5 \times 10^4$ J/m$^3$. b) Time evolution of spin wave emission from a domain wall at 9 GHz after the onset of the ac spin-polarized current. The domain wall is located at $x = 0$ and the dotted line indicates the location where $m_z$ drops below $0.15 \times 10^{-3} M_s$.

position where the spin wave amplitude is drops below $0.15 \times 10^{-3} M_s$. After approximately 15 excitation periods, the system reaches a steady-state motion and the energy from the excitation signal is in balance with the internal spin wave damping. Due to the symmetry of the system, the wavefront of the propagating spin waves is oriented parallel to the domain wall, i.e., the angle between magnetization and the wave-vector is $45^\circ$.

Similar simulations were performed at other frequencies. By spatial Fourier transformation of a steady-state spin wave profiles, a full dispersion diagram was obtained (Fig. 6.2b). In the dispersion graph, the horizontal line at 9.3 GHz corresponds to the non-dispersive domain wall resonance mode and the additional dispersive modes at $f_{res}/n$ corresponds to higher harmonics ($n$ is an integer). Above a frequency of 24 GHz ($\lambda \sim 150$ nm), the domain wall can not follow the microwave current and the emission of spin waves is strongly suppressed.

The emission of spin waves from domain walls in multiferroic heterostructures or other anisotropy-modulated magnetic films can be exploited in devices where multiple spin wave sources are required and the position of the excitation point is important. An example of such a system is illustrated in Fig. 6.3a. Here, two nanowires are placed on top of a ferroelectric boundary. Due to the perfectly straight nature of the ferroelectric domain wall, the excitation point is identical in both nanowires (the accuracy of the relative positions is not limited by the resolution of lithographic process). Furthermore, our simulations indicate that the wavelength of the emitted spin waves in this constrained geometry can be reduced significantly due to the formation of Damon-Eshbach-like standing spin
Results

Figure 6.3. a) A schematic illustration of two parallel nanowires placed on top of a ferroelectric domain boundary with 90° rotation of polarization. The double headed arrows indicate the direction of uniaxial magnetic anisotropy induced by strain transfer from the substrate. b) Simulated dispersion relation of spin waves radiated from the domain boundary. c) Snapshots of the spin wave profiles at different resonance frequencies of the domain wall oscillator.

Wave modes inside the domain wall. Profiles of emitted spin waves for a 400 nm wide magnetic nanowire are shown in Fig. 6.3c and their dispersion is depicted in Fig. 6.3b. For high frequencies, the domain wall modes converge causing highly monochromatic excitation of spin waves. At 50 GHz, for example, the wavelength of emitted spin waves is 25 nm and the wavelength distribution is only 1 nm.

6.2 Spin wave confinement in anisotropy modulated systems

Figure 6.4. a) Magneto-optical Kerr microscopy image of the CoFeB/BaTiO₃ sample at zero magnetic field. The arrows denote the direction of magnetization inside the domains. b) Polar plot of the remanent magnetization in two neighbouring stripes. c) Easy and hard-axis hysteresis loops of two neighbouring stripe domains.
In publication V, we studied spin wave dynamics in a BaTiO$_3$/50 nm CoFeB/5-
Au nm sample with the same ferroelectric domain configuration as described
in the previous section. The magneto-static properties were investigated using
magneto-optical Kerr microscopy (Fig. 6.4). From these static measurements, a
highly regular pattern of magnetic domain walls in the CoFeB film was inferred
and the widths of the two domain types were estimated as 5 $\mu$m and 7.5 $\mu$m.
Successful imprinting of uniaxial magnetic anisotropy was also confirmed by
measurements of the remanent magnetization in two neighbouring stripe do-
main. Fig. 6.4b shows a polar plot of the local remanent magnetization as
a function of magnetic field angle, demonstrating 90° rotation of the uniaxial
anisotropy axis at the domain boundary. From the hard-axis hysteresis loops
(Fig. 5.3c), we derived an anisotropy strength of 1.4 × 10$^4$ J/m$^3$.

![Figure 6.5. a) Broadband spin wave spectroscopy measurement geometry (not to scale). The
domain boundaries are oriented parallel to the signal line of the coplanar waveguide. b) Experimentally measured angular dependence of spin wave excitations in the
CoFeB/BaTiO$_3$ sample at 60 mT in-plane magnetic field. First order ($n=1$) and
higher-order confined spin wave modes ($n=1\ldots9$) are indicated c) Micromagnetic
simulations of the spin wave excitation spectrum as a function of magnetic field
angle. In the simulation, alternating stripe domains with a width of 5 and 7.5 $\mu$m
and other experimental parameters are assumed.](image)

The magnetodynamic properties of the sample were investigated using all-
electrical broadband spin wave spectroscopy. In the experiment, the sample
was placed face-down on a coplanar waveguide with a 20-$\mu$m-wide signal line
that produces a long-wavelength microwave field with a main excitation at
$k \approx 0.096$ rad/$\mu$m. The corresponding wavelength of this excitation field amounts
to about 65 $\mu$m. Since this is about 10 times the width of the individual stripe
domains, it can be regarded as a reasonably uniform excitation field. The ex-
perimental geometry is shown in Fig. 6.5a. During spectroscopy measurements,
a 60 mT strong external magnetic field was applied along different in-plane
directions of the CoFeB film in 2.5° steps (Fig. 6.5b). This field is more than twice
the saturation field of the hard-axis hysteresis loop (Fig. 6.4c) and, thus, guar-
antees a nearly uniform magnetization state without magnetic domain walls.
The measurements condition in this study are thus very different from those
in the previous section where spin wave emission from pinned domain walls in
the remanent state of the magnetic film was considered. If the field is applied
parallel ($\theta=0^\circ$) or perpendicular ($\theta=90^\circ$) to the magnetic stripe domains, the
resonance frequencies are the same in both domains and the magnetization
precesses uniformly. Rotation of the magnetic field aligns the magnetization

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more parallel to the easy and hard axis of the a1-a2 domains. This modifies the relative strengths of their effective magnetic fields and leads to a separation of resonance frequencies. In addition to the two main branches, additional modes are measured between these resonances (Fig. 6.5b). These modes are attributed to higher order confined modes (up-to the 9th order) in stripes with lower effective field.

![Figure 6.6](image)

To elucidate the origin of the spin wave confinement, we performed micromagnetic simulations using MuMax3. The anisotropy modulated system was simulated using four stripe domains with alternating widths of 5 and 7.5 μm and periodic boundary conditions were applied in the film plane to mimic an infinite film. The simulation mesh was divided into $3.05 \times 3.05 \times 50$ nm$^3$ cells and the anisotropy axis was rotated abruptly at the domain boundaries. The experimentally determined anisotropy constant of $1.4 \times 10^4$ J/m$^3$ was used together with an exchange constant of $2.1 \times 10^{-11}$ J/m, a Gilbert damping constant of 0.005 and a saturation magnetization of $1.15 \times 10^6$ A/m. The saturation magnetization was measured using vibrating sample magnetometry on a 50 nm CoFeB reference film grown on a SrTiO$_3$ substrate. The spin wave modes were excited by applying a 0.1 mT spatially uniform Gaussian magnetic field pulse perpendicular to the stripe domains. After excitation, the z-component of magnetization was averaged over the simulation area and recorded as a function of time for 10 ns. Resonance spectra were obtained from these data by Fourier transformation of the out-of-plane magnetization component ($m_z$) as function of the external magnetic field angle. The simulated angular dependence of the spin wave spectrum is shown in Fig. 6.5c. The simulations reproduce the main features, including the excitation of higher-order confined modes for angles 0-45° and 135-180°, as in the experimental spectrum (Fig. 6.5b).

To confirm the formation of standing spin waves, the dispersion relations were calculated for the two stripes when a 60 mT field was applied at 135°, i.e., parallel to the hard axis of the wider 7.5 μm stripe (Fig. 6.6b). In the stripe with higher effective field, no modes exist below 10 GHz. In contrast, modes exist...
in the neighbouring stripes where the magnetization is aligned parallel to the hard axis. This mode discrepancy allows formation of intense confined modes up-to ninth order whereafter modes exist in both areas. At frequencies above this limit, emission of spin waves from the boundary reduces the intensity of the standing spin waves. Emission of spin waves is discussed in detail in the next sections. To illustrate the spin wave profile, the system was simulated using a continuous ac-magnetic field at the frequency of standing spin wave modes. After steady-state precessional motion was reached, the spatial distribution of the z-component of magnetization was recorded for two nanoseconds. As an example, the spin wave profile of the $n = 7$ mode is shown in Fig. 6.6c.

### 6.3 Spin wave excitation from a standing spin wave mode

The anisotropy modulated system discussed in the previous section can also be used to emit short-wavelength spin waves from the boundaries by exciting standing spin wave modes in stripe domains with large effective field. This process takes place at high frequencies where a confined mode in one of the domains and a propagating mode in neighboring domains coexist (Fig. 6.7a). The wavevector of the propagating spin wave is much larger that that of the standing spin wave (see arrow in Fig. 6.7a). To illustrate this phenomenon, additional simulations were performed with 22.5 $\mu$m and 7.5 $\mu$m wide stripe domains. In these simulations, the magnetic bias field was applied parallel to the easy anisotropy axis of the narrow domain. The other domain with small effective field was intentionally widened to avoid formation of an interference pattern. The location of the first and third order spin-wave modes were determined first by a Gaussian field pulse and these modes were excited subsequently using an uniform ac magnetic field. After the onset of excitation, a standing spin wave is formed in the narrow domain, i.e., the domain with large effective magnetic field (Fig. 6.7b). In addition to the buildup of the confined mode, a propagating spin wave is excited in the neighboring domain from the anisotropy boundary.
The emitted spin wave has considerable amplitude and shorter wavelength than the confined mode. This mode conversion can be explained by downwards shift of the dispersion relation at the domain boundary. At the standing spin wave frequency of the stripe with larger effective field, a short wavelength mode exist in the domain with small effective field. This conversion process is illustrated in Fig. 6.7a by the arrow. Thus, the domain with large effective field acts as a local spin wave source. In contrast, the perfectly confined mode discussed in the previous section does not excite propagating spin waves. In the standing spin wave frequencies of the stripe with small effective field, the wave vector becomes imaginary at the stripe boundary. Since no modes exist in the neighboring domains, propagation into these areas is forbidden and the z-component of magnetization decays quickly (Fig. 6.6c).

The wavelength of emitted spin waves can be tuned by rotating the external magnetic field. Rotation of the bias field changes the effective field and the angle between magnetization and the spin wave vector. Rotation of the field away from the respective hard and easy anisotropy axis in two neighboring domains moves the dispersion curves in Fig. 6.7a towards each other. In addition, if the field rotates more perpendicular to the exciting CPW (i.e. towards a backward-volume spin-wave configuration), the slopes of the two dispersion curves decrease. The latter effect significantly decreases the wavelength of the emitted spin waves (Fig. 6.7c). Using experimentally derived parameters, we simulate a minimal wavelength of 108 nm for a magnetic field angle of 72°.

### 6.4 Spin wave emission from an isolated anisotropy boundary

![Figure 6.8](image)

**Figure 6.8.** a,b) Simulated $m_z$ component of magnetization at two different time steps after steady-state precessional motion is reached. The excitation frequency is 40 GHz. The profiles are recorded with a phase difference of about 75°. c) Simulated (symbols) and calculated (lines) dispersion relations for simulated (symbols) and theoretical model (line) of emitted spin waves from an isolated anisotropy boundary. The red curve corresponds to spin waves propagating in the domain where the field is parallel to the easy anisotropy axis (large effective magnetic field) and the blue curve for the neighboring domain where the field is perpendicular to the easy anisotropy axis (small effective magnetic field).

Davies and Kruglyak demonstrated that spin waves with wavelengths down to several hundreds nanometers can be emitted from an isolated boundary of a
Results

semi-infinite film by spatially uniform excitation field [27]. In this geometry, the emission of unidirectional spin waves is triggered by a dynamic demagnetization field that pins spins at the edge of the structure. In our BaTiO$_3$/CoFeB system, we find that a uniform microwave magnetic field also triggers spin wave emission from an isolated anisotropy boundary. In our geometry, the emission of spin waves is bi-directional. Moreover, rather than a dynamic demagnetization field, the emission of spin waves from anisotropy boundaries is caused by the excitation of dissimilar forced magnetization precessions in neighboring domains with different effective magnetic fields. The thus formed time-dependant divergence of magnetization at the boundary produces a dynamic exchange torque which launches spin waves into the domains. This excitation mechanism is illustrated by the simulations in Fig. 6.8a,b showing the spin wave profiles of propagating spin waves at 40 GHz for two different phases of the excitation signal. The simulations indicate a steady-state precessional motion for a 60 mT field parallel to the easy and hard anisotropy axis of neighboring domains. The two snap shots illustrate how the amplitudes of magnetization precession at the anisotropy boundary vary over time (Fig. 6.8b). As a result, a dynamic exchange torque is generated which emits spin waves. Figure 6.8c shows the dispersion relations of emitted spin waves in the two domains with different effective magnetic field. The shape of the dispersion curves illustrates how the character of emitted spin waves changes gradually from magnetostatic (large wavelength) to exchange-dominated (short wavelength) with increasing frequency. As a guide for the eye, the theoretical dispersion relation of the system was calculated using Eq. 3.9.

6.5 Programmable spin-wave filtering by pinned magnetic domain walls

Spin-wave based logic devices require a method to control spin-wave transmission. A well-studied technique is based on reconfigurable magnonic crystals where the lateral modulation of magnetic properties can be actively turned on or off [104, 18, 115]. Such systems have been realized by utilization of Oersted fields in a periodic conductor array [18]. An alternative method relies on the creation of periodic heating gradients in a YIG film by high-intensity laser pulses [112]. Magnonic crystals require a few repetitions of the modulation period to efficiently block spin-wave transmission via the formation of forbidden bands. In publication VII, we investigate spin-wave filtering by a firmly pinned 90° domain walls of a multiferroic 50 nm CoFeB/BaTiO$_3$ bilayer (the same structure as described in Sections 5.2-5.4). In zero magnetic field, two stable domain wall states are possible: one with a 90° head-to-tail magnetic structure and another in which the magnetization aligns in a 90° head-to-head or tail-to-tail configuration [37]. The widths of these two types of domain walls are very different. The width of the broad head-to-head (or tail-to-tail) domain wall increases rapidly with film thickness due its dependence on magnetostatic energy (Publication I
and III). In contrast, the width of the head-to-tail configuration is determined by a competition between exchange and anisotropy energies and, hence, it only weakly depends on film thickness. In the experiments on spin wave filtering, the CoFeB film thickness is 50 nm which ensures full imprinting of the stripe domain pattern and allows tuning of the domain wall width from 50 nm to 1.6 μm by non-volatile switching between the two wall types.

![Image](image_url)

**Figure 6.9.** a) Schematic illustration of the experimental configuration. Spin waves are excited using a microwave antenna on top of the CoFeB film and the propagating spin waves are probed using μ-BLS across the wall (orange arrow). The lower panel displays magneto-optical Kerr microscopy image of the area. b) BLS measurement measured 500 nm from the antenna. c) Phase-resolved μ-BLS scans across a head-to-head (black) and head-to-tail (red) domain walls.

The magnetostatic properties of the 28 nm TaO_x/3 nm Ta/50 nm CoFeB/BaTiO_3 sample were investigated using magneto-optical Kerr microscopy. The measurements confirmed successful imprinting of the anisotropy pattern with the anisotropy axis rotating by 90° at the domain boundaries. The oxide layer on top of the CoFeB film was introduced to avoid shorting of the antenna structures that were patterned on top. The antennas consisted of 3 nm Ta/50 nm Au, they were 500 nm wide, and were aligned parallel to the pinned magnetic domain walls (Fig. 6.9a). The positions of the domain walls with respect to the antennas were measured using magneto-optical-Kerr microscopy. In the spin-wave transmission experiments described below, the antenna edge and domain-wall center are separated by approximately 2 μm.

The transmission of propagating spin waves through pinned magnetic domain walls was investigated using phase-resolved micro-focused Brillouin light scattering (μ-BLS). First, a BLS frequency sweep was measured 500 nm from the antenna. In this experiment, the frequency of the signal source was swept between 7 and 13 GHz in discrete steps with sub-second dwell time (Fig. 6.9b). The BLS signal was recorded for 3000 scans to locate the excitation and sensitivity maxima of the sample. It was observed that the spin-wave excitation efficiency from the antenna was highly non-reciprocal. The signal strength was reduced by one order of magnitude when the direction of magnetization was reversed underneath the antenna. To maintain consistency of the experimental measurement between the two domain wall states and to maximize the experimental sensitivity, the magnetization underneath the antenna was aligned in the same
way for the head-to-tail and head-to-head configurations. The largest sensitivity in the BLS frequency sweep was attained at 10.85 GHz and, therefore, this frequency was used in subsequent μ-BLS scans across the pinned magnetic domain wall. Figure 6.9c shows spin-wave transmission results for a narrow head-to-tail (red curve) and broad head-to-head (black curve) domain wall. In the head-to-head configuration, spin waves with frequency of 10.85 GHz and a wavelength of approximately 1.3 μm are transmitted through the domain wall. In contrast, almost no spin-wave signal was observed behind the head-to-tail domain wall.

To further analyze spin-wave transmission through 90° magnetic domain walls, we performed micromagnetic simulations in MuMax3. In the simulations, we considered two 20-μm-wide stripe domains. Two dimensional periodic boundary conditions were applied in the CoFeB film plane to mimic an infinite film and the cell size was set to 2.4×9.6×12.5 nm³. The anisotropy pattern was imprinted by abrupt 90° rotation of the anisotropy axis at the domain boundary. Spin waves were excited locally by an out-of-plane sinusoidal magnetic field at the center of one of the domains, i.e., 10 μm from the pinned domain wall. The system was first driven for 14 ns which is sufficient to reach a steady-state precessional motion and subsequently the $m_z$ component of magnetization was recorded for one nanosecond in 2.5 ps steps. A weak Bloch nature of the domain wall was removed by averaging the $m_z$ along the thickness direction. Results of the simulations for a frequency of 10.85 GHz are shown in Fig. 6.10. The simulations reproduce the large difference in spin-wave transmission between narrow head-to-tail and broad head-to-head domain walls, as observed in μ-BLS experiments.
In the broad domain wall, the effective magnetic field reduces slowly inside the wall and the magnetization rotates more perpendicular to the wave vector of the propagating wave. As a result, the dispersion relation is shifted downwards and the magnetization rotates towards the Damon-Eshbach configuration. Rotation towards the Damon-Eshbach configuration increases the slope of the dispersion relation and, as a consequence, the group velocity is enhanced. These changes are illustrated in Fig. 6.10a,c) where the wavelength and the decay of the propagating spin wave is reduced inside the domain wall. The enhancement of the group velocity is illustrated in Fig. 6.10a where the dashed line shows the envelope function of the spin waves without a magnetic domain wall. Thus, the head-to-head domain walls is fully transparent to spin waves. In contrast, waves are strongly reflected by the narrow head-to-tail domain wall (Fig. 6.10d,f). The effective magnetic field peaks at the domain wall and the peak is surrounded by two field minima. This non-uniform field generates a resonance mode at the domain wall that is characterized by oscillation of two out-of-phase antinodes on the opposite sides of the wall (inset in Fig. 6.10f). Spin waves are absorbed and reflected by this resonance mode and this produces an interference pattern. Interference is illustrated by the undulations of spin-wave maxima and minima in the contour graph (Fig. 6.10f).

The ability to reversibly turn spin wave transmission on and off by reversible non-volatile switching between two distinct domain wall states could be exploited in magnon conduits, filters, logic gates or memory elements. Based on our results on spin-wave transmission through single 90° domain walls, we propose a new structure for active spin-wave manipulation. The device concept consists of three
stripe domains with two pinned domain walls. In this system, the magnetization reversal in the center domain switches the domain wall configuration between head-to-head/tail-to-tail and head-to-tail/head-to-tail. Switching between these states drastically alters the spin wave transport through the domain walls at the domain wall resonance frequency. This effect is shown in Fig. 6.11 for a frequency of 11 GHz. In the simulations, the pinned domain walls were separated by 1.5 μm which is sufficient for stabilizing a head-to-head or tail-to-tail domain walls. The broad domain walls are fully transparent to propagating spin waves (Fig. 6.11a,b). In contrast, the excitation of the domain wall resonance mode in the head-to-tail configuration reduces the transmission to nearly 0% (Fig. 6.11c,d). Since the spin-wave signal can be easily turned on or off by magnetic switching of one single domain, we refer to our structure as a magnetic spin-wave valve.

6.6 Spin wave emission via a dynamic interface exchange torque in YIG/CoFeB bilayers

Magnonics opens new prospects for data processing, storage and transmission of information using the amplitude or phase of spin waves [59, 98, 51, 22, 84]. Ferrimagnetic YIG has captivated large interest due to its small Gilbert damping factor. YIG has been utilized in various spin-wave components such as transistors [16], magnonic crystals [21, 17], logic gates [95] and multiplexers [26]. In these experiments, propagating spin waves are usually excited by microwave antennas. However, the wavelength is limited by the excitation structure and the small saturation magnetization of YIG allows only spin-wave frequencies in the low GHz at small magnetic bias fields. Larger wave vectors and frequencies can be attained by using grating couplers at the expense of excitation efficiency [118, 119]. Alternatively, perpendicular standing spin waves (PSSW) can be excited in magnetic films. The wave vector of these modes is approximately \( k = p\pi/d \) where \( p \) is the order number and \( d \) is the film thickness. In nanometer thick films, the wave vector of PSSW modes is large and their frequency is high. The formation of PSSWs can be induced, for example, by spatially non-uniform excitation from a microwave antenna [52], asymmetric spin pinning at the interfaces of a film [118] or by abrupt heating using intense laser pulses [91]. In publication VIII, we studied the emission of high-frequency spin waves and the formation of perpendicular standing spin waves in exchange coupled YIG/CoFeB bilayers using uniform microwave field excitation. Our results reveal a new excitation mechanism that is driven by a dynamic exchange torque at the interface.

In our study, epitaxial 295 nm thick YIG films were grown on (111) oriented GGG substrates using pulsed laser deposition and good epitaxial growth was confirmed by XRD. The pure YIG films were characterized by all-electrical broadband spin wave spectroscopy using a flip-chip technique where the sample
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Figure 6.12. a) Schematic illustration of the measurement geometry. b) Linescan at 20 mT for pure 295 nm YIG (green) and a 295 nm YIG/50 nm CoFeB bilayer (orange). The perpendicular standing spin waves in YIG are labeled with $p = 1, 2, \ldots$ and the ferromagnetic resonance modes with $p = 0$. c) Experimental spin wave spectrum as function of magnetic field for a YIG/CoFeB bilayer. d) Spin wave spectrum at high magnetic fields.

was placed face-down onto a coplanar waveguide with a 50 $\mu$m-wide signal line and 800 $\mu$m-wide ground lines. The gap between the signal and ground lines was 30 $\mu$m. The CPW was chosen to excite spin-waves with $k \approx 0$. Consequently, excitation of perpendicular standing spin waves via a spatially non-uniform microwave field is avoided. The measurement geometry is shown in Fig. 6.12a. Next, a 50 nm CoFeB and thin Ta capping layer were grown using magnetron sputtering on top of the same YIG films. The spin wave absorption spectra was recorded using all-electrical broadband spectroscopy as function of magnetic field. The bias field was applied along the coplanar waveguide. The results for pure 295 nm YIG film and 295 nm YIG/50 nm CoFeB are compared in Fig. 6.12b. The linescan measured at 20 mT on the pure YIG films shows an intense peak at about 2 GHz, which corresponds to uniform precessional motion, and perpendicular standing spin wave modes are absent due to near uniform microwave excitation field. In contrast, the measurement on the YIG/CoFeB bilayer shows multiple modes in addition to the ferromagnetic resonance modes (FMR) in YIG and CoFeB. These modes correspond to higher order PSSWs in the YIG film and their field dependence is shown in Fig. 6.12c,d. The PSSW modes are especially strong near the FMR resonances of YIG and CoFeB, as illustrated
by experiments at high magnetic fields (Fig. 6.12d) where modes up to tenth order are efficiently excited. The excitation mechanism is based on a dynamic exchange torque at the interface. Near the FMR frequencies, the difference in forced magnetization precession in the YIG and CoFeB films is maximized and, via interface coupling, this produces a strong dynamic exchange torque. Strong coupling between the layers produces characteristic anti-crossings and mode hybridization.

Figure 6.13. a) Simulated (top panel) and experimental (bottom panel) spin-wave spectra for 295 nm YIG/50 nm CoFeB (orange) and 295 nm YIG/10 nm Ta/50 nm CoFeB (dashed green) at a magnetic bias field of 30 mT. b,c) Simulated spin wave spectra as function of magnetic field for YIG/CoFeB and YIG/Ta/CoFeB, respectively.

To elucidate the origin of PSSW excitation further, the YIG/CoFeB bilayer was modeled using micromagnetic simulations. In the simulations, we used experimental parameters for the saturation magnetization and exchange constants and a relatively large Gilbert damping constant of 0.005. The damping parameter was artificially increased to reduce the simulation time but this did not affect the results. The spin waves were excited using a spatially uniform 3 mT sinc-type magnetic field pulse with a cut-off frequency of 20 GHz. Interface exchange coupling is determined in MuMax by calculating the harmonic mean value of the exchange interaction field. Simulations for a 295 nm YIG/50 nm CoFeB bilayer are shown in Fig. 6.13 and they are compared to an experimental spin wave spectrum. The resonance frequencies and line-shapes of the peaks are in good agreement with each other. However, the amplitudes of the peaks differ significantly because the parameters that are extracted from experiments and simulation are different. In the experiments, the intensity of the signal is determined by inductive coupling between the magnetization and the microwave current. Consequently, the absorption increases at higher frequencies if the precessional angle remains constant. In contrast, the intensity in the simulated spectra is proportional to the precessional amplitude and its strength is
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independent of frequency.

Figure 6.13c also shows simulations for a 295 nm YIG/10 nm Ta/50 nm CoFeB trilayer. Just like experiments on this structure (see Publication VIII), the simulations demonstrate that PSSWs are not excited in the YIG film when exchange-coupling at the YIG/CoFeB interface is turned off by insertion of a thin Ta layer.

Figure 6.14. a,b) Spatial distribution of spin-wave modes in 295 nm YIG/50 nm CoFeB and 295 nm YIG/10 nm Ta/50 nm CoFeB. c,d) Time evolution of the z-component of magnetization at 4.9 GHz for the same structures. The excitation frequency corresponds to the $p = 4$ PSSW mode.

Simulations of the spatial distribution of excited spin-wave modes are shown in Fig. 6.14a,b) for the 295 nm YIG/50 nm CoFeB bilayer and the 295 nm YIG/10 nm Ta/50 nm CoFeB trilayer. The results are obtained by a Fourier imaging technique where the time evolution of $m_z$ was Fourier-transformed on a cell-by-cell basis. The spatial simulations illustrate the formation of symmetric and strongly confined modes in the YIG film. If the coupling between the layers is suppressed by a non-magnetic 10 nm Ta layer, only uniform precessional modes are simulated.

PSSW modes are excited by constructive interference of spin waves that are emitted from the YIG/CoFeB interface if one of the perpendicular confinement conditions is fulfilled. The time evolutions for the YIG/CoFeB and YIG/Ta/CoFeB are depicted in Fig. 6.14c,d). The simulations illustrate the magnetization response after the onset of uniform ac-field excitation at $t = 0$ s. After the
microwave field is turned on, propagating spin waves are launched from the interface towards the YIG film at a wavelength of about 150 nm. This wavelength matches the $p = 4$ PSSW condition and therefore, it is resonantly enhanced (Fig. 6.14c). In contrast, if exchange coupling between the CoFeB and YIG layers is turned off by a non-magnetic Ta spacer layer, only forced magnetization precessions with dissimilar amplitudes are excited (Fig. 6.14d).

![Figure 6.15. a,b) Simulated amplitude of $m_z$ and $m_x$ in a 295 nm YIG film (blue lines) and a 50 nm CoFeB layer (orange line) as a function of frequency. Besides the FMR maxima, the magnetization of both films is forced to precess over a large frequency range.](image)

We explain the emission of spin waves from the YIG/CoFeB interface by a dynamic exchange torque originating from non-equal forced precessions in the YIG and CoFeB layers. Forced magnetization precessions are excited in both layers over a large frequency range, as illustrated by the simulations of Fig. 6.15. However, the difference of the forced precessional motion is never zero and, as a consequence, spin waves are excited over a broad frequency range. When the wavelength of the excited spin wave satisfies the PSSW confinement condition, the mode is resonantly amplified. The difference in forced oscillations is maximized near the FMR modes of the CoFeB and YIG films and, as a result, the dynamic exchange torque is largest in their vicinity. This enhancement of excitation efficiently is clearly seen in experiments (Fig. 6.12). While the intensities of the PSSW modes are large near FMR resonances, they reduce between these modes because the amplitudes of forced magnetization precessions become more similar (Fig. 6.15).
7. Conclusions

In this thesis, I presented results on the excitation of confined and propagating spin waves in multiferroic heterostructures and YIG/CoFeB bilayers. The spin waves are either excited by an uniform microwave magnetic field or spin-polarized electric current. The investigated multiferroic heterostructures are based on BaTiO$_3$/CoFe and BaTiO$_3$/CoFeB bilayers with regular rotations of uniaxial magnetic anisotropy. In zero magnetic field, the anisotropy boundaries strongly pin 90° magnetic domain walls. The microstructure of the domain walls can be reversible switched between two types: broad head-to-head or tail-to-tail and a narrow head-to-tail walls. When the narrow domain wall is driven into oscillation using a spin-polarized current, its back-and-forth motion emits spin waves into the domains. The wavelength of the spin waves can be scaled down to sub-100 nm when the ferromagnetic film is patterned into nanowires. The same anisotropy boundary can be used to coherently excite spin waves with a narrow wavelength distribution in parallel nanowires. Alternatively, anisotropy-modulated multiferroic heterostructures can be exploited to confine higher order spin-wave modes and excite propagating spin-waves in structurally continuous ferromagnetic films. This behavior, which is experimentally demonstrated in magnetic saturation, is explained by a modulation of the effective magnetic field. When a uniform microwave magnetic field is applied, perfectly confined higher-order standing spin waves are excited in domains with smaller effective field. At higher frequencies, the confined spin wave modes are also excited in the domains with larger effective field. In this case, propagating spin waves are emitted from the boundary into neighboring areas with smaller effective field. The wavelength of the propagating spin waves is smaller in comparison to those of the standing spin wave modes. When the magnetization is rotated by an external bias field, the wavelength of emitted spin waves is continuously tuned. Furthermore, I demonstrated that a single magnetic anisotropy boundary in a continuous magnetic film can be used to excite short-wavelength spin waves by an uniform microwave field at any frequency. This spin-wave excitation mechanism is explained by dissimilar forced magnetization precessions in neighboring anisotropy domains. These forced precessions produce a time-varying divergence of magnetization at the boundary, which is counteracted
Conclusions

by bi-directional launching of spin-waves. I also showed efficient excitation of perpendicular standing spin waves in YIG/CoFeB bilayers based on a similar excitation mechanism. The large difference in the saturation magnetization generates dissimilar forced precessional amplitudes between the layers when they are excited using an uniform microwave field. As a result, a dynamic exchange-torque is generated at the interface which launches spin waves into the YIG film. A perpendicular standing spin wave (PSSW) is formed when the wavelength of emitted spin waves matches the confinement condition. In the experiments, PSSW modes were measured up to tenth order. The excitation efficiency is maximized near the ferromagnetic resonance frequency of the CoFeB or YIG film. In these frequency ranges, the difference in forced precessional amplitudes is maximized and the generated dynamic exchange torque is largest. Finally, I proposed a spin-wave valve based on 90° Néel domain walls in a BaTiO3/CoFeB multiferroic heterostructure. In zero magnetic field, domain walls are strongly pinned onto the anisotropy boundaries and their structure can be reversibly switched between a broad head-to-head (or tail-to-tail) and narrow head-to-tail configurations. The broad domain walls are fully transparent to propagating spin waves while spin-wave transmission is reduced to nearly zero in the narrow head-to-tail configuration. A system consisting of two parallel pinned domain walls could be exploited as non-volatile spin-wave memory or in magnonic logic devices. Here, switching of the magnetization in the center domain toggles the configuration of the system from two narrow to two broad domain walls, or vice versa. As a result, spin wave transmission is turned on or off at will in a non-volatile manner.
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