LIGHT SCATTERING STUDIES OF METALLIC MAGNETIC MICROSTRUCTURES

DISSERTATION

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By

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ABSTRACT

In this thesis, the physics underlying the magnetic behavior of metallic microstructures, including their responses to magnetic fields and electric currents is explored. The dynamic and static components of the magnetization are respectively probed through Brillouin light scattering and Kerr imaging method. The design, growth and fabrication of various structures are presented, while the experimental findings are analyzed by theoretical modeling and calculations. The highlights include (a) Brillouin light scattering studies of spin precession under tunable magnetic field imbalance, (b) Kerr imaging of layer-by-layer magnetic reversal in cobalt-platinum multilayer, (c) Observation of spin-polarized current induced domain wall motion in magnetic microwires. All of these results demonstrate that light scattering as an excellent tool for probing novel functionality of metallic magnetic microstructures. Future prospects along the direction of research involved in this thesis are also presented.
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TABLE OF CONTENTS

Abstract...........................................................................................................ii

Acknowledgments..........................................................................................iii

Vita................................................................................................................vi

List of Tables.................................................................................................x

List of Figures...............................................................................................xi

Chapters:

1. Introduction...................................................................................................1

2. Characterization techniques..........................................................................4

   2.1 Brillouin light Scattering...........................................................................4
   2.1.1 Scattering profile.................................................................................4
   2.1.2 Experimental setup of BLS.................................................................6
   2.1.3 Polarization of photons scattered from magnons.........................8
   2.1.4 Scattering intensity..........................................................................10
   2.1.5 New trends of BLS.................................................................14
   2.1.6 Comparison with other spin wave detection techniques.........15

2.2 Magneto-optical Kerr effect......................................................................16

   2.2.1 MOKE magnetometry.....................................................................20
   2.2.2 Time averaging.................................................................................25
   2.2.3 The magnitude of Kerr rotation....................................................27
   2.2.4 Kerr imaging....................................................................................29
   2.2.5 Comparison of MOKE with other magnetometry techniques.....30
   2.2.6 New trend in MOKE magnetometry and microscopy..............31

3. Micro-fabrication techniques.......................................................................33

   3.1 Film deposition.....................................................................................33
      3.1.1 Substrate.......................................................................................33

    vii
# LIST OF TABLES

<table>
<thead>
<tr>
<th>TABLE</th>
<th>PAGE</th>
</tr>
</thead>
<tbody>
<tr>
<td>3.1</td>
<td>36</td>
</tr>
</tbody>
</table>

Comparison of properties between different substrates
<table>
<thead>
<tr>
<th>FIGURE</th>
<th>PAGES</th>
</tr>
</thead>
<tbody>
<tr>
<td>2.1</td>
<td>(a) Scattering of a laser photon by a bulk magnon. (b) Scattering of a photon by a surface magnon. (c) Incident laser beam, magnetization and spin wave wavevector…………………………………………………………………..…….6</td>
</tr>
<tr>
<td>2.2</td>
<td>Experimental setup of BLS……………………………………………………..…8</td>
</tr>
<tr>
<td>2.3</td>
<td>Focusing of a laser beam by a lens………………………………………………15</td>
</tr>
<tr>
<td>2.4</td>
<td>Three types of MOKE scattering geometry(a) Longitudinal MOKE (b) Polar MOKE (c) Transverse MOKE …………………………………………………..19</td>
</tr>
<tr>
<td>2.5</td>
<td>(a) Experimental setup of MOKE magnetometry (b) Analyzing polarizer and Kerr rotation of the optical electric field…………………………………………………………………..21</td>
</tr>
<tr>
<td>2.6</td>
<td>Contribution of each term in equation (2.10) to intensity $I$ (arbitrary unit) at photodiode as polarizer analyzing angle $\delta$ varies (in degrees). The absolute magnitudes of each term are displayed in (a) while the associated noise due to a fluctuation factor $F$ (equals 1.01) is presented in (b)…………………………….24</td>
</tr>
<tr>
<td>2.7</td>
<td>(a) Time averaging acts as a low pass or band pass filter to selectively pick up the signal. MB stands for measurement bandwidth. (b) Tilted background of hysteresis loop arises due to Faraday rotation in camera lens. (c) Hysteresis loop not returning back to the original position (d) A more ideal hysteresis loop……26</td>
</tr>
<tr>
<td>2.8</td>
<td>Kerr imaging system set up during the present investigation.........................30</td>
</tr>
<tr>
<td>2.9</td>
<td>(a) Kerr image of $50 \times 100 \mu$m channel without mirror rotation (b) With mirror rotation…………………………………………………………………..32</td>
</tr>
<tr>
<td>3.1</td>
<td>Schematic of a thermal evaporator………………………………………………..37</td>
</tr>
<tr>
<td>3.2</td>
<td>Schematic of a diode sputtering machine…………………………………………39</td>
</tr>
<tr>
<td>3.3</td>
<td>Illustration of magnetron sputtering……………………………………………41</td>
</tr>
<tr>
<td>3.4</td>
<td>Diffusion pump……………………………………………………………………42</td>
</tr>
</tbody>
</table>
3.5 Picture of blades of a turbomolecular pump……………………………………..43

3.6 Steps of photolithography with subtractive techniques (a) Deposition (b) Spinning of photoresist (c) Exposure (d) Development (e) Etching and (f) Stripping off photoresist………………………………………………………………………..47

3.7 (a) Contact lithography (b) Projection lithography (c) Olympus Optical Microscope BH-2. The arrow indicates the slot to place in a mask for projection lithography……………………………………………………………………….49

3.8 Steps of lithography with additive techniques (a) Spinning of photoresist (b) Photolithography (c) Deposition of the magnetic films (d) Lift-off………………….51

3.9 Schematic illustration of electron beam lithography………………………………54

3.10 AFM image of a 2 micron wide, 50 nm thick permalloy wire fabricated on the surface of SiO$_2$/Si………………………………………………………………...54

4.1 Schematic sketch of experimental setup…………………………………………57

4.2 Brillouin spectra recorded at fixed external field ($H_{ext}$) of 600Oe as a function of current…………………………………………………………………………....59

4.3 Schematic illustration of in-phase (SA) and out-of-phase (SO) magnetization oscillations shown on precession cones for various Amperian ($H_{cur}$) and external field ($H_{ext}$) directions and magnitude in case of $|H_{ext}| > |H_{cur}|$ (parallel alignment of magnetization). The rectangles represent the two magnetic layers in cross-section. The probe laser beam is incident from the left and directions of $I_{channel}$ indicated by ⊗ and ⊘…………………………………………………………………………..61

4.4 Calculated mode profiles of SO mode (solid lines) and SA mode (dotted lines) for different free current density $J_f$ flowing in the Cu spacer layer………………………………………………………………………………66

4.5 Brillouin spectra recorded with a fixed external field $H_{ext} = 50$Oe as a function of $I_{channel}$………………………………………………………………………..68

4.6 Schematic illustration of in-phase (SA) and out-of-phase (SO) magnetization oscillations shown respectively on precession cones for various Amperian ($H_{cur}$)
and external field ($H_{\text{ext}}$) directions and magnitude in case of $|H_{\text{ext}}| < |H_{\text{cur}}|$ (anti-parallel alignment of magnetization). The rectangles represent the two magnetic layers in cross-section. The probe laser beam is incident from the left and directions of $I_{\text{channel}}$ indicated by ⊗ and ◊.

4.7 Same as Fig. 4.3 except the external field $H_{\text{ext}} = 500\text{eOe}$, that is smaller than the current induced field $H_{\text{cur}} = 1000\text{eOe}$. The direction of the overall magnetization $M_o$ is thus primarily determined by $H_{\text{cur}}$.

4.8 Anti-Stokes Brillouin spectra recorded in sample #2 for various values of $H_{\text{ext}}$ and $H_{\text{cur}} (I_{\text{channel}})$.

4.9 Solid lines represent calculated switching strength parameter $S$ as a function of $|H_{\text{ext}}|$ for various values of $|H_{\text{cur}}|$. An experimental point is hollow (filled) if spectra correspond to anti-parallel (parallel) alignment of magnetization is detected.

4.10 (a) Coupled pendulums with different local gravitational fields ($g + \Delta g$) and ($g - \Delta g$). (b) and (c) illustrate dependence of normalized eigenvector elements $\rho_1$ and $\rho_2$ on $\eta (= \Delta g/\alpha)$ for the lower (l) and higher (h) frequency normal modes.

5.1 Illustration of spin-torque transfer mechanism.

5.2 Image of channel utilized (a) A wide field view of the structure. The central channel is connected to large area of film on one side, and by a series of micro-wires on the other side. (b) A high magnification view of the central channel. The dimension of the channel width/length = 8/125 $\mu$m.

5.3 Experimental setup for applying electrical pulse through our fabricated sample.

5.4 Response (Kerr images) of a head to head domain wall to electrical current pulses.

5.5 Response (Kerr image) of a tail to tail domain wall to electrical current pulses.

5.6 Example of the response of a tail-to-tail domain wall to electrical current from another sample with identical pattern dimension.
5.7 (a) Spins are aligned along the x axis. A domain wall is placed at the center. Exchange interaction is turned off ($A=0$). (b) x component of the monitored spin. (c) y component of the monitored spin (d) z component of the monitored spin. Only the relative (not absolute) values of the amplitude components matter but not their absolute value. Current density is set at $14 \times 10^{11} \text{Am}^{-2}$. (e) x component of the monitored spin against time if a current density of $20 \times 10^{11} \text{Am}^{-2}$ is used instead. The unit for time is picosecond. Since vertical axis in (b) to (e) corresponds to spin x component oscillation amplitude, they are unit-less.

5.8 Amplitude of the monitored spin (a) x component. (b) y component and (c) z component. Initially, a current of $14 \times 10^{11} \text{Am}^{-2}$ is applied with the electron flow along -x and exchange stiffness $A=0$. At $t=170$ ps in the figure, A abruptly increases to $0.3 \text{pJ/m}$. Time t is in picoseconds.

5.9 Coordinate system of the magnetic moment m. L is the grid size utilized in the simulation.

5.10 Result of micromagnetic simulation. (a) Vortex domain wall (b) Transverse domain wall.

6.1 (a) Room temperature coercivity ($H_C$) for [Co(4 Å)/Pt($t_{Pt}$)]$_N$ multilayers with repetition $N = 5$. (b) Hysteresis loop for sample with $N = 5$, $t_{Pt} = 11$ Å (arrow in (a)). (c) – (g): Some Kerr microscopy images of [Co(4 Å)/Pt(11 Å)]$_5$ recorded at consecutive number of field pulses at $H_p$ indicated in (b) created by current pulse flowing through field coils. (i) to (l): Kerr images with pulses $H_p$ applied in the reversed direction. (h) and (m): Kerr images of the multilayer after saturated by a large magnetic field opposite to the initial magnetization.

6.2 (a) Room temperature coercivity ($H_C$) for [Co(4 Å)/Pt($t_{Pt}$)]$_N$ multilayers with repetition $N = 5$. (b) Hysteresis loops for sample with $N = 5$, $t_{Pt} = 41$ Å (arrow in (a)). (c) – (g): Some Kerr microscopy images of [Co(4 Å)/Pt(41 Å)]$_5$ recorded at consecutive number of field pulses at $H_p$ indicated in (b). (i) to (l): Kerr images with pulses $H_p$ applied in the reversed direction. (h) and (m): Kerr images of the multilayer after saturated by a large magnetic field opposite to the initial magnetization.

6.3 (a) Room temperature coercivity ($H_C$) for [Co(4 Å)/Pt($t_{Pt}$)]$_N$ multilayers with repetition $N = 8$. (b) Hysteresis loops for sample with $N = 8$, $t_{Pt} = 41$ Å (arrow in (a)). (c) – (g): Some Kerr microscopy images of [Co(4 Å)/Pt(41 Å)]$_8$ recorded at consecutive number of field pulses at $H_p$ indicated in (b). (i) to (l): Kerr images
with pulses $H_p$ applied in the reversed direction. (h) and (m): Kerr images of the multilayer after saturated by a large magnetic field opposite to the initial magnetization………………………………………………………………………………115

6.4 (a) Room temperature coercivity ($H_C$) for [Co(4 Å)/Pt($t_{Pt}$)]$_N$ multilayers with repetition N = 12. (b) Hysteresis loops for sample with N = 12, $t_{Pt} = 41$ Å (arrow in (a)). (c) – (i) Some Kerr microscopy images of [((4 Å)/Pt(41 Å))]$_{12}$ consecutive number of field pulses at $H_p$ indicated in (b)……………………………………117

6.5 Kerr images of magnetic reversal for [Co(4 Å)/Pt(23 Å)]$_5$ multilayers……….119

6.6 Kerr images of magnetic reversal for [Co(4 Å)/Pt(33 Å)]$_5$ multilayers………120

6.7 Kerr images of magnetic reversal for [Co(4 Å)/Pt(41 Å)]$_{12}$ multilayers………122

6.8 Kerr images of magnetic reversal for [Co(4 Å)/Pt(41 Å)]$_{12}$ multilayers………123
CHAPTER 1

INTRODUCTION

Magnetism has contributed in important ways to human beings. The utilization of magnetic phenomena dates back to five thousand years where the legendary emperor of China, named Huang di, won a battle over his enemy by the usage of a “South-pointing vehicle”, essentially a magnetic compass. For centuries, magnetism has assisted humans to navigate the ocean and conquer the world. In the 19th Century, the discovery of transformation between electricity and magnetism by Michael Faraday brought new applications of magnetism such as electric generators that has revolutionized the modern world. In the middle of the last century, magnetism found another important application: as recording media in audio and video cassette and memory elements in computers and calculators. Since then, properties of magnetic microstructures have attracted much attention as improvement of these recording and memory utilities required deep understanding of the underlying physics. The discovery of the giant magnetoresistance effect [1] has led to ten fold of increase in memory density on computer hard disks over the last decade [2].

In present day physics, the study of magnetic memory systems is integrated into a broader category named “spintronics” [3], where in the “spin” (the basic unit of 1
magnetism) is not only utilized as a memory element but also gives rise to computation capabilities. The field of spintronics is usually sub-divided into two broad divisions depending on the material system on which they are based: (a) semiconductors [4] and (b) magnetic metallic structures [5]. The interest in manipulating the spin in semiconductors mainly lies in the hope of integration with the mature semiconductor industry. Semiconductors provide “gain” (amplification of signal), an important ingredient which is lacking in their metallic counterparts. However, magnetic metallic materials remain an equally important topic in spintronics, largely due to their convenience in memory applications. “Magnetic Random Access Memory” (MRAM), which mainly utilize magnetic metallic materials, has already entered the era of commercialization and is expected to replace the traditional dynamic random access memory (DRAM) in computer chips [5]. Moreover, there exist creative approaches for using magnetic metallic microstructures to perform computation and logic operation [6,7]. When combined with microelectromechanical system (MEMS), metallic magnetic materials lead to novel imaging applications of magnetic resonance force microscopy (MRFM) [8]. Such metallic structures are being applied to process microwave signals [9] and even as a microwave source themselves [10].

From the above examples, the versatility of application of magnetic metallic microstructures is evident. Characterization of magnetic metallic microstructures is usually through magnetic force microscopy (MFM) [11] and giant magnetoresistance (GMR) phenomenon [12]. This thesis concentrates on the study of magnetic metallic structures using optical methods. The two main optical methods: Brillouin light scattering (BLS) and Magneto-optical Kerr effect (MOKE) are introduced in Chapter 2. These
two methods are respectively adopted to probe two novel physical phenomena in magnetic metallic microstructures. By using BLS, the modification of spin-precession in a magnetic double layer by the Amperian magnetic field associated with an electrical current traversing through the middle spacer metallic layer was discovered. Such a phenomenon cannot be easily detected using traditional techniques such as MFM and GMR, since they are largely insensitive to the high frequency dynamics of the spins in the structures. These results are discussed in Chapter 4. On the other hand, by setting up the MOKE imaging system in the laboratory, we have pursued two projects: First, we provide experimental observation of domain wall motion in magnetic metallic wires directly induced by spin-polarized current. These experimental results together with theoretical analysis are discussed in Chapter 5. Second the first evidence of layer by layer magnetic reversal of Co/Pt multilayers, which exhibit perpendicular magnetic anisotropy, at calibrated Co layer separation has been provided. Such films have great potential for high density information storage in the future. These discoveries are presented in Chapter 6. A summary of the thesis and possible future studies are discussed in Chapter 7.
CHAPTER 2

CHARACTERIZATION TECHNIQUES

2.1 Brillouin light Scattering

Brillouin light scattering (BLS) had emerged as an important tool in studies of magnetic metallic single and multi-layers in the 1980s and 1990s [13-15]. The underlying principle is based on inelastic scattering of laser light from spin waves residing in the magnetic layers. It involves the emission or absorption of a magnon (particle counterpart of spin wave) by a photon while the photon is inelastically scattered within the sample medium. These two scattering processes are named “Stokes” and “Anti-Stokes” process respectively.

2.1.1 Scattering profile

Figure 2.1(a) schematically illustrates the scattering of a photon by a bulk magnon. The out-coming photon must have a wave-vector magnitude very close to that of the in-coming photon in order to satisfy energy conservation, since the energy of the magnon absorbed or emitted is much less than the photon itself. Therefore the out-coming photon wave-vector k’ must lie on the dash line circle, which represents the sphere of photon wave-vector with energy equal to that of the in-coming photon. During the experiment,
a camera lens is placed in proximity to the sample. The cone of photons collected in Figure 2.1(a) represents the range of scattered photon wavevectors collected by the camera lens. The collection geometry in Figure 2.1(a) is named “backscattering”, as the cone of photon collection has its central axis aligned with the incoming photon wavevector.

Figure 2.1(b) illustrates the scattering of a photon by a surface magnon. Since surface spin waves have a decaying amplitude profile (instead of a propagating profile) in a direction normal to the opaque film, surface magnons observed by light scattering have their momentum undefined perpendicular to the surface. Therefore momentum is only to be conserved in plane along the sample surface. The vertical dash line in Figure 2.1(b) act as a guide to the eye, illustrating that the horizontal momentum component is conserved after scattering. It is interesting to note that the wave-vector and energy exchange between the photon and the surface magnon depends on the angle of incidence of the incoming light, while no such dependence exists for scattering of bulk magnons (Assuming that $k$ and $k'$ are collinear, i.e. $k'=-k$, the wavevector magnitude of the emitted or absorbed bulk magnon $|q|$ always equal 2$|k|$, while for surface magnon $|q|$ equal 2$|k|\sin\theta$, where $\theta$ is the angle between $k$ and sample surface normal). Therefore, by varying the incident angle of the photons, signal from bulk magnon reveals no new information but the surface magnon shows the important $\omega$-$k$ (energy-momentum) dispersion relation. Important parameters such as exchange-stiffness constant can be subsequently deduced [13].
2.1.2 Experimental Setup of BLS

The experimental BLS setup of is depicted in Figure 2.2. Photons from an Argon-
ion laser source are split into two unequal parts by a beam splitter. The major part goes to
the sample and the resulting scattered light is directed into the interferometer after
traversing various optical elements. The second, low intensity, laser beam is named the
“reference beam”. It enters the interferometer without interaction with the sample. The
system requires this reference beam to keep the Fabry-Perot mirrors in alignment. The
two light beams are directed to pass through the Fabry-Perot mirrors six times before
entering the detector. By passing the beam multiple times through the mirrors, the
contrast between stop and pass frequency is increased five to six orders of magnitude
when compared with using a single pair of mirror. Suppose $C$ is the transmission contrast
ratio for a single mirror pair, where $C = T_1/T_2$ with $T_1$ and $T_2$ the maximum and minimum
transmission values defined in the inset of Fig. 2.2, then multi-passing the laser beam $n$
times through the mirror corresponds to an effective transmission ratio of $C^n$ [13]. During
the experiment, the spacing within the two Fabry-Perot mirrors are continuously varied
simultaneously with a period of roughly one second to obtain the required analysis of
photon frequency distribution of the scattered photons. For detailed workings of the two
Fabry-Perot mirrors refer to Ref. 13.
Figure 2.2: Experimental setup of BLS. M: Mirror; S: Sample; L: Lens; CP: Prism; BS: Beam Splitter; S1: Shutter 1; S2: Shutter 2; R: Reflecting prism; FP: Fabry-Perot; P: Pinhole; D: Detector; Inset (upper right corner): Transmission $T$ of light against frequency $f$ through a single pair of Fabry-Perot mirrors. $T_1$ and $T_2$ identify the maximum and minimum transmission value.

### 2.1.3 Polarization of photons scattered from magnons

Light consists of a fluctuating electric field. When such an oscillatory electric field $\mathbf{E}$ is present in a magnetic material with magnetization $\mathbf{M}$, an oscillatory electric polarization $\mathbf{P}$ will be produced according to the formula (in S. I. units) [14]

$$\mathbf{P} = (\varepsilon_{11} - 1)\mathbf{E} + K\varepsilon_0 (\mathbf{E} \times \mathbf{M}) / M_s$$  \hfill (2.1)
Where \( M \) is the magnitude of \( \mathbf{M} \), \( \varepsilon_{11} \) is the diagonal component of the dielectric constant matrix of the magnetic material, and \( K \) is the magneto-optic coefficient. Values of \( \varepsilon_{11} \) and \( K \) are frequency dependent and their values at two common spectroscopy laser lines (514.5nm and 632.8nm) for Ni, Co and Fe can be found in Ref. 14. Suppose \( \mathbf{M} \) also contains a fluctuating component, then \( \mathbf{M} \) can be written as \( \mathbf{M} = \mathbf{M}_0 + \mathbf{m} \), where \( \mathbf{M}_0 \) is the static component and \( \mathbf{m} \) is the oscillatory component. It readily follows from equation (2.1) that \( \mathbf{P} \) will contain a term proportional to \( \mathbf{E} \times \mathbf{m} \).

Since in thin films it is the surface spin wave, and not the bulk magnon, which carries information on the dispersion relation, we focus our study on the surface spin wave. For surface spin waves, the magnon travels largely around the direction perpendicular to the magnetization (for details see Ref. 13). Therefore, in order to probe the surface spin wave, the incident plane of the laser beam should lie parallel to the magnon wavevector, as shown in Figure 2.1(c). The laser polarization is defined as “p” (“s”) polarized if the oscillating electric field lies within (perpendicular to) the incident plane. In Figure 2.1(c), the incoming light is p-polarized, and its electric field is given by \( \mathbf{E} = E_x \mathbf{e}_x + E_y \mathbf{e}_y \). Its cross-product with \( \mathbf{m} = m_x \mathbf{e}_x + m_y \mathbf{e}_y \) results in an electric dipole moment pointing in the \( z \) direction: \( \mathbf{P} = (E_x m_y - E_y m_x) \mathbf{e}_z \). The resulting oscillating electric dipole radiates electromagnetic waves with a characteristic spatial pattern [16]. An array of electric dipoles residing on the sample surface and oscillating in the \( z \)-direction will ensure the radiated electromagnetic wave has its electric field along the \( z \) direction. Therefore the scattered light is s-polarized, which is perpendicular to p-polarized incident light.
Suppose instead the incident light is s-polarized, i.e. the electric field is given by
\[ \mathbf{E} = E_z \mathbf{e}_z . \] We can verify that the resultant electric polarization is proportional to
\[ E_z (m_x \mathbf{e}_y - m_y \mathbf{e}_x) . \] Therefore in this case the scattered light is p-polarized. Thus we conclude that the polarization of light scattered by a magnon is always perpendicular to the polarization of the original incident light. This differs from light scattered by acoustic phonon that has its polarization direction unchanged. With analyzing polarizer selected to proper orientation, we can selectively transmit the spin wave signal to pass through while the entire phonon signal is being blocked. The efficiency of the polarizer is revealed in the Brillouin spectra shown in chapter 4. While these spectra clearly display spinwave signals, no trace of phonon peaks are evident confirming the integrity of the polarizer.

2.1.4 Scattering intensity

We are now ready to provide a rough estimate of the BLS intensity. From equation (2.1) we know that the oscillatory polarization due to magnetization fluctuation has a magnitude of
\[ P = Ke_0 E \theta_m \] (small angle approximation is used for \( \theta_m \)), where \( \theta_m \) is the precession angle of the magnetization. The radiation power of a point electric dipole with oscillation amplitude \( p \) equals [16]
\[
\text{Power} = \frac{\mu_0 p^2 \omega^4}{12 \pi c} \tag{2.2}
\]
Let us assume the sample surface is composed of an array of these oscillating electric dipoles with a lateral separation $r$. Let $d$ be the light penetration depth. Then,

$$p = PV = K\varepsilon_0 E \theta_m r^2 d$$  \hspace{1cm} (2.3)

If $q$ is the wavevector of the spin wave of interest, then

$$q = \frac{4\pi}{\lambda} \cos \theta$$  \hspace{1cm} (2.4)

where $\lambda$ and $\theta$ are the wavelength and incident angle of the laser beam respectively. We assume the separation of the oscillating electric dipole mentioned above approximately equals half of the spin wave wavelength, i.e. $r = \pi / q$ (since two spinwave with equal amplitude and wavelength traveling oppositely produce a standing wave with a periodicity equal to half the wavelength). If we suppose the laser beam spot covers a square of area $R \times R$ on the sample surface, then the number of radiating dipoles will equal $(R/r)^2$. On the other hand, the electric field strength $E$ associated with the laser light is related to the incident laser power $P_{in}$ through:

$$P_{in} = \frac{1}{2} c \varepsilon_0 E^2 R^2$$  \hspace{1cm} (2.5)

In seeking for an expression for the magnetization precession angle $\theta_m$, the energy density stored by a magnetization $\mathbf{M}$ in a magnetic field $\mathbf{H}$ is given by $E = -\mu_0 \mathbf{H} \cdot \mathbf{M}$. If $\mathbf{M}$ deviates from the direction of $\mathbf{H}$ by a small angle $\theta_m$, then it can be easily verified that

$$E = -\mu_0 H M_s \left(1 - \theta_m^2\right)^{1/2} \approx -\mu_0 H M_s \left(1 - \left(1/2\right)\theta_m^2\right),$$

therefore the energy density stored in this spin wave mode equals $(1/2)\mu_0 H M_s \theta_m^2$. Equating this energy to $k_B T / (R^2 t)$ (we
assume the spin wave is confined within the $R \times R$ square area), where $t$ is the magnetic film thickness, we can calculate $\theta_m$. Finally, by substituting the expression for $p$ of (2.3) into equation (2.2) and take account of the number of radiating dipoles and collection angle, the radiated signal power is given by

$$P_{\text{signal}} = \frac{\mu_0 \theta_m^4}{12 \pi c} \left( K \varepsilon_0 E \theta_m r^2 d \right)^2 \Delta \Omega \left( \frac{R}{r} \right)^2$$

(2.6)

where $\Delta \Omega$ is the stereoradian angle of photon collection. The signal photon flux is thus given by

$$\text{Flux} = \frac{P_{\text{signal}}}{hf}$$

(2.7)

Including the following numerical values: $R=50 \text{ } \mu \text{m}$, $P_m=100 \text{ } \text{mW}$, $c=3 \times 10^8 \text{ } \text{m}^{-1}\text{s}^{-1}$, $\varepsilon_0=8.85 \times 10^{-12} \text{ } \text{F/m}$, $\lambda=514.5 \text{ } \text{nm}$, $\mu_0=1.26 \times 10^{-6} \text{ } \text{H/m}$, $K=\sqrt{0.15}$ [14], $d=5 \text{ } \text{nm}$, $\Delta \Omega = (4/5)^2 / (4\pi)$ (assuming collection camera lens is 5 cm from the sample and has a diameter of 4 cm), $\theta=45^\circ$, $t=30 \text{ } \text{nm}$, $H=600 / (4\pi \times 10^{-3}) \text{ } \text{Am}^{-1}$, $M=8 \times 10^5 \text{ } \text{Am}^{-1}$, $k_B=1.38 \times 10^{-23} \text{ } \text{J/K}$, $T=300 \text{ } \text{K}$. We can verify that the signal flux equals $\sim 4400$ photons per second. This is the contribution from one spin wave mode. Suppose a free spectra range of 30GHz is being utilized in the BLS interferometer and there are approximately 1000 channels, then the width of each BLS channel equals 30GHz/1000=30MHz. As mentioned, the spin wave are confined in an area with lateral dimension of $R=50 \text{ } \mu \text{m}$. Given that the velocity of spin wave is typically 1000 ms$^{-1}$, we can verify that the frequency spacing of spin wave mode on the magnetic sample roughly equals 60 MHz. Therefore each BLS channel is sensitive to approximately half a spin wave mode and
hence the theoretical photon counts in each channel after 1000 passes (1 pass=1ms) should be 4400 divided by 2 equals 2200. Experimentally, each channel only receives about 100 counts after 1000 passes. Therefore there are 2 orders of magnitude difference between the estimate and experiment. According to Ref. 14, components in the optical path (mirrors, lens etc) and the BLS interferometer can lead to attenuation of signal by a factor of 20, while another factor of 10 comes from the quantum efficiency of the photodetector. The remaining discrepancy may come from the fact that we have neglected the attenuation of laser light electric field inside the magnetic film, the modification of the electric field by the substrate as well as the surface roughness. Given the uncertainty in all the parameters involved, the discrepancy between the estimated and measured of BLS spin wave signal is reasonable.

During the BLS experiment, the elastic scattered beam from the sample is blocked since we are only interested in the inelastic scattered light signal. However, the system requires a reference beam directly from the laser to keep the Fabry-Perot mirrors aligned. The power of the reference beam is required to be approximately 1mW to enable the feedback mirror aligning system to work properly. By using numerical values estimated above, we can calculate the ratio of spin wave signal power to the reference beam power to be roughly $10^{-12}$ while the multi-pass interferometer has a transmission contrast greater than $10^9$ [13], spin wave signal and reference beam power at the detector are roughly similar, which produce the largest limiting factor in our detection sensitivity. For the light detector installed in our laboratory, the dark count is ~5 per second (i.e. 50 per 1000 passes), which is much smaller to the strength of our spin wave signal (see above). Therefore with the estimated spin wave signal strength, in order to overcome the
noise embedded in the leakage of reference beam into the system, the spectrum must be collected for several thousand passes. Therefore it takes approximately an hour to collect a BLS spin wave spectrum with acceptable signal to noise ratio.

2.1.5 New trends of BLS

While the basic design of the BLS apparatus has essentially remained unchanged for two decades [13], new features are being developed to upgrade the system. Two major extensions have been realized in recent years: (a) time resolved BLS [17] and (b) micro-Brillouin measurement [18]. In the time resolved work, a pulse generator is used to initiate a spin wave soliton inside a magnetic sample. The same electric pulse is used to trigger the light detection system with a controllable time delay. By varying this time delay, the time evolution of the spatial profile of the spin wave soliton can be tracked. The time span for the whole evolution process usually lasts for several hundred nanoseconds and the time resolution is reduced to several nanoseconds. The response time of electronic components is one of the major factors limiting the time resolution.

In a micro-Brillouin experiment, the relatively large camera lens of a conventional set up is replaced by a microscope objective, which has a much shorter focus length. The focusing of a laser beam of wavelength $\lambda$ by a lens is shown in Figure 2.3. The minimum width of the laser beam, called the beam waist ($w$), is related to the converging angle $\theta$ by the formula

$$w = \frac{2\lambda}{\sin(\theta/2)} \quad (2.8)$$

Therefore the smaller the focal length of the lens, the larger is $\theta$ and the smaller $w$. 

14
and beam size on the sample surface. In a conventional BLS experiment, the laser spot diameter on the sample is roughly 50 µm. However, in a micro-Brillouin experiment, the spot size on the sample is limited by diffraction to several hundred nanometers. This significantly improves the spatial resolution of the BLS detection. However, since the signal intensity is proportional to the interaction area of the laser with spin wave, the diminished spot size decreases the signal strength sequentially if the areal power density remains the same. Hence such micro-Brillouin experiments are generally limited to detection of artificially excited spin wave (by means of external microwave source), while detection of room temperature, thermally activated, spin wave is yet to be realized.

![Figure 2.3: Focusing of a laser beam by a lens.](image)

2.1.6 Comparison with other spin wave detection techniques

Traditionally, spin waves in magnetic materials are detected using neutron scattering [19]. In neutron scattering, the momentum of the neutron is typically much larger than the photons in a BLS experiment. By varying their momentum, neutron scattering is able to probe spin wave over the entire Brillouin zone. This contrasts
BLS which is limited only to spin waves near zone center. However, due to the small scattering cross-section from magnons, neutron scattering can generally probe only spin waves in bulk material and detection of spin wave in thin films has been largely unsuccessful. Recent positive attempts to push the limit of neutron scattering in thin films have been reported [20].

New generation of synchrotron sources have been installed in national laboratories of various countries. They provide opportunity of inelastic x-ray scattering experiments. Since x-rays also have much larger momentum than optical photons used in BLS, x-ray scattering can probe the entire Brillouin zone. One successful attempt has been made to detect phonons with x-rays [21]. However, the study of magnons by x-ray scattering has still not been possible due to the weakness of scattering signal.

Ferromagnetic resonance (FMR) [22] is an important traditional tool to probe spin wave excitations. The line-width of the FMR resonance signal reveals information on spin wave damping, while the line-width of a BLS signal is largely controlled by instrumental broadening and yields no information on damping. However, FMR generally is only sensitive to the zone center \( k=0 \) spin wave (uniform mode). The recently developed of ferromagnetic resonance force microscopy (FMRFM) [23] is a new approach that has higher sensitivity. While pulsed inductive microwave magnetometer (PIMM) [24] was built following similar ideas as FMRFM, BLS will retain its unique status in probing spin waves in magnetic metallic multilayers into the foreseeable future.

### 2.2 Magneto-optical Kerr effect

The rotation of light polarization when transmitted through or reflected from
a magnetic material is known as the “Faraday effect” or “Kerr effect” respectively. Ref. 25 provides a very good account on the discovery of these two phenomena. Within this thesis, we are dealing mainly with magnetic metallic materials and, due to their metallic character, are largely opaque. Therefore the “Kerr effect” is utilized. The magneto-optical Kerr effect is often referred by the acronym MOKE.

While BLS exploits the dynamic, fluctuating component of $\mathbf{M}$ in equation (2.1), MOKE basically takes advantage of the static part. This leads to a MOKE scattering geometry which is different from BLS. Three types of MOKE scattering geometry are displayed in Figure 2.4, namely (a) longitudinal MOKE, (b) polar MOKE and (c) Transverse MOKE. In longitudinal MOKE, the sample magnetization lies in the sample surface plane, which is the same as in BLS. However, in contrast to BLS, the plane of incidence lies parallel, instead of perpendicular, to the direction of magnetization. In Figure 2.4(a), the polarization of the incident light is aligned along the $y$ direction (s-polarized). Since the magnetization is along $x$, by the $\mathbf{E} \times \mathbf{M}$ term in equation (2.1) results in an oscillating electric field along the $z$ direction. Therefore the reflected light contains a p-polarized component. The magnitude of this p component has been exaggerated in Figure 2.4(a) as its magnitude is much smaller than the s-polarized component in the reflected light. It can be easily verified that if the incident light is p-polarized, a new s-polarized component will be generated in the reflected light. Therefore, incident light of any polarization will have its polarization “rotated” upon reflection and the direction of this rotation depends on the orientation of the sample magnetization. We can thus gain knowledge of the sample magnetization direction by evaluating the polarization rotation of the reflected light. One important property of longitudinal
MOKE is that light must be incident \textit{obliquely} on the sample. In Figure 2.4(a), if light incident normal to the sample surface, the reflected light will travel normal to the surface (i.e. in z direction). The newly generated oscillating polarization and therefore the electric field, originated form the term $\mathbf{E} \times \mathbf{M}$ in equation (2.1) will also lie in z-direction. Since light is a transverse wave and doesn’t oscillate along its direction of propagation, the reflected light will not pick up this newly generated oscillating electric field. Therefore the reflected light undergoes no rotation. In summary: Longitudinal MOKE is absent when the light incident is perpendicular to the sample surface.

Due to strong demagnetization effects by the shape anisotropy, the magnetization of a magnetic thin film usually lies in the sample surface plane. However, some magnetic films possess strong crystal anisotropy such that the magnetization of these films naturally lies perpendicular to the surface. This is illustrated in Figure 2.4(b) where the magnetization is aligned along z. If the polarization of the incident light is along the y, then according to equation 2.1, an oscillating electric field along x-direction will be generated in the reflected light. This scattering geometry is referred to as the “polar” MOKE. Note that, in contrast to the longitudinal MOKE, the signal strength of the polar MOKE is strongest when the light is incident perpendicular to the sample surface. Since in polar MOKE geometry, the reflected light electric field oscillates exactly in the direction of $\mathbf{E} \times \mathbf{M}$ ($\mathbf{E}$ is the electric field of the incident light), while in longitudinal MOKE the same condition doesn’t hold, therefore polar MOKE signal is usually stronger than longitudinal MOKE.

Figure 2.4(c) depicts the transverse MOKE scattering geometry. In this geometry the magnetization, lies perpendicular to the incident plane. Suppose the incident
light is s-polarized (aligned along $y$) direction, then the incident light polarization lies in the same direction as the magnetization and $\mathbf{E} \times \mathbf{M}$ equals zero, i.e. results in no rotation. If the incident light is p-polarized, we can easily verify that the newly generated oscillating electric field also lies in the incident plane and yields no rotation of the polarization. Therefore the $\mathbf{E} \times \mathbf{M}$ term in equation (2.1) doesn’t give rise to any polarization rotation and any MOKE signal that arises will be due to higher order effects. Ref. 26 reports a successful magnetometry study on bulk iron by the transverse MOKE.

Figure 2.4: Three types of MOKE scattering geometry. (a) Longitudinal MOKE (b) Polar MOKE (c) Transverse MOKE
2.2.1 MOKE magnetometry

One of the main objectives of magnetometry is to obtain the hysteresis loop of the sample magnetization against an external magnetic field. The experimental setup of a longitudinal MOKE magnetometry experiment is displayed in Figure 2.5 (a). A narrow parallel laser beam is converged to a focused spot on the sample. Just prior to the sample the light is vertically polarized. The reflected light is transmitted through an analyzing polarizer (the “analyzer”) before being converged onto a silicon photodiode by a camera lens. It is important to place the polarizers immediately before and after the sample. No optical components, not even the converging and collecting lenses are placed between the polarizer and the sample to ensure any polarization rotation associated with light intensity changes at the detector is solely from the sample. Thus, stray magnetic fields do not lead to Faraday rotation of the light beam within the glass material lenses and contribute to light intensity changes at the photodiode. The polarizer must be selected carefully. A Glan-Thompson type polarizer under the influence of stray magnetic fields can give rise to undesirable Faraday rotation. A sheet polarizer is preferred and it should withstand the laser power during the experiment (~80mW within a spot of diameter ~1mm). The mirror in front of the CCD camera is removable. This allows the CCD camera to monitor the position of the laser spot on the sample surface.
Figure 2.5: (a) Experimental setup of MOKE magnetometry (b) Analyzing polarizer and Kerr rotation of the optical electric field. The two dotted lines represent the horizontal and vertical direction. The solid lines represent the polarization of light which is allowed to transmit through the analyzer, which is tilted away from absolute horizontal with an angle $\delta$.

Figure 2.5 (b) depicts the orientation of the analyzer ($\delta$) and the Kerr rotation ($E_p/E_s$) of the laser beam. Light is vertically polarized by the polarizer immediately prior to the sample, and is therefore s-polarized light strikes the sample ($E_s$). After reflection the beam gains a p-polarized component ($E_p$) and enters the analyzer. The axis of
transmission polarization of the analyzer is tilted at an angle $\delta$ to the horizontal. The light intensity beyond the polarizer is given by

$$I = \left| E_s \sin \delta + E_p \cos \delta \right|^2$$  \hspace{1cm} (2.9)

Suppose $E_s = |E_s|$ and $E_p = |E_p|e^{j\phi_p}$, after some manipulation and noting that the light is depolarized, namely by factor of $D$ (where $D << 1$), we obtain

$$I = \left| E_s \right|^2 \sin^2 \delta + \left| E_p \right|^2 \cos^2 \delta + \left| E_s \right|\left| E_p \right| \cos \phi_p \sin(2\delta) + DE_s^2$$  \hspace{1cm} (2.10)

In Figure 2.6(a), the contribution of each term in equation (2.10) to the intensity $I$ is plotted as a function of the analyzer orientation angle $\delta$. In this figure, we adopt the following numerical values: $E_s = 1.0$ (arbitrary unit), $E_p = E_s (50 \times 10^{-3}) \pi / 180$ (50 milli-degree of Kerr rotation), and $\phi_p = 0$ (no Kerr ellipticity for simplicity). The term $DE_s^2$ reflects the depolarization effect, which arise due to depolarize scattering happens at the sample and polarizer. Assuming an extinction ratio of 10000:1 for the polarizer, we use $D = 0.0001$. Both the second (quadratic term) and third term (dc term) of equation (2.10) contains information on the Kerr rotation ($E_p$). However from Figure 2.6(a) it follows that when $\delta$ is not equal to zero, the dc term provides a much stronger signal than the quadratic term. Therefore MOKE magnetometry should be carried out with $\delta$ slightly deviated from the extinction condition to take advantage of the dc term of equation (2.10). Although the dc term grows with $\delta$, so does the first term of (2.10) (see Figure 2.6(a)). Therefore $\delta$ cannot be too large or the detector will be saturated by the contribution from the first term.
Figure 2.6(b) investigates the choice of $\delta$ from the point of view of noise. Since the power of the laser and so $E_s$ fluctuates, there will be noise arising from the first term of equation (2.10). The same happens with the depolarizing term. They are given by

\[
\left( |FE_s^2 - E_s^2| \right)^2 \sin^2 \delta \quad \text{and} \quad \left( |FE_s^2 - E_s^2| D \right)^2
\]

respectively. For a 1% laser power fluctuation, $F$ equals 1.01. It is evident from Figure 2.6(b) that the noise associated with the first term of (2.10) grow swiftly with $\delta$, and the best signal to noise occurs around 0.5 degree.
Figure 2.6: Contribution of each term in equation (2.10) to intensity $I$ (arbitrary unit) at photodiode as polarizer analyzing angle $\delta$ varies (in degrees). The absolute magnitudes of each term are displayed in (a) while the associated noise due to a fluctuation factor $F$ (equals 1.01) is presented in (b). Note in (a), the weakness of $E_p^2 \cos^2 \delta$ leads to its overlap with the x axis.
2.2.2 Time Averaging

As seen from Figure 2.6 (b), even when $\delta$ is optimized to value around 0.5°, the signal to noise ratio (SNR) lies between 3:1 to 4:1. Time averaging is needed to enhance the SNR. The noise associated with the laser fluctuation is mainly from short term ripples. On the other hand, the Kerr signal is produced by slowly sweeping the external magnetic field (for a single loop). Therefore the short term ripples distribute at relatively high frequency when compared to the MOKE signal for a single loop (see Figure 2.7a). By time averaging (taking signal over a period of time and averaging), we simply introduce a low pass filter into our measurement system and discard the high frequency ripples associate with the laser and electronics. Figure 2.7 (b) displays a problematic hysteresis loop. The loop is tilted and opened. While the sharp transition near the center represents Kerr rotation signal from the sample, the tilted background is from Faraday rotation occurring in the converging and collecting lens. As previously noted, by placing the polarizer between the sample and the lens, this tilt problem can be remedied. In Figure 2.7 (c), the tilt is suppressed but the loop is not closed. This figure was created by spending time on signal averaging in obtaining each data point of the loop, while the entire loop is sampled once only. Time averaging in this manner will eliminate short terms ripple of the MOKE system, but the loops are still susceptible to long term drift within the setup. In contrast to short term ripples which happen at the sub-second level, long term instabilities usually takes minutes to develop and lead to failure in obtaining a close loop. Figure 2.7 (d) represents an ideal hysteresis loop. It is neither tilted nor open. The important point is that instead of signal averaging to obtain each data point of the loop and sampling the entire loop only once, it is better to obtain multiple complete
loops (ten or more) and take their average. In this way, both the short term ripples and the long term drift will be suppressed while retaining the same time to complete the entire scan. Since in this multi-loops scan circumstance, each loop is completed quickly, the signal is being actually modulated at high frequency (see Fig. 2.7a) providing, in essence, a lock-in detection technique.

Figure 2.7 (a) Time averaging acts as a low pass or band pass filter to selectively pick up the signal. MB stands for measurement bandwidth. (b) Tilted background of hysteresis loop arises due to Faraday rotation in camera lens. (c) Hysteresis loop not returning back to the original position (d) A more ideal hysteresis loop.
2.2.3 The magnitude of Kerr rotation

The magnitude of the Kerr rotation can be estimated following the line of thought in section 2.1.4 for the BLS intensity. The only change is setting $\theta_m$ (angle between $\mathbf{M}$ and $\mathbf{H}$) to equal $\pi$ (In MOKE, the static component of magnetization reverses direction and therefore twice the magnitude of the magnetization contributes to the signal, instead of only the fluctuating component as in the BLS). For the numerical values suggested in section 2.1.4, the magnitude of the Kerr rotation is given by $\sqrt{P_{\text{signal}}/P_{\text{in}}}$ equals 300 milli-degrees. This value is approximately ten times higher than the measured value. As mentioned (section 2.1.4), many factors including neglecting the optical electric field attenuation inside the sample will lead to overestimation of the Kerr signal. Ref. 27 provides a more rigorous calculation of the Kerr signal in various magnetic films, which takes complex dielectric constant tensor of the material into account.

2.2.4 Kerr Imaging

In MOKE magnetometry, the signal is obtained by averaging over the entire area covered by the laser spot on the sample surface (roughly a diameter of 50µm). The magnetization profile may however not be homogeneous within the focused beam. In order to spatially resolve the local magnetization, Kerr imaging is required.

Figure 2.8 displays the Kerr imaging system. It was found that a home-made microscope constructed of different optical components was not able to achieve the require stability. Physical vibrations that occur in such a home-made microscope were unacceptable at the highest magnifications. The vibration problem was suppressed when a commercial microscope was used instead.
In Figure 2.8, the light first passes through a Glan-Thompson polarizer. The polarizer orientation can be rotated to tune the incoming laser polarization. The laser beam is then reflected by a rotating mirror. The surface plane of the mirror is slightly off from being absolutely perpendicular to its rotational axis. In this way, when the mirror rotates, the light beam is reflected in slightly different directions, as shown in Figure 2.8. With the help of subsequent lenses and mirrors, the light beams are then converged back to the same spot on the sample surface.

The reason for introducing the rotating mirror is to eliminate interference effects: Figure 2.9(a) displays an image of a 50 × 100 micron channel on the sample surface obtained without the rotating mirror and significant interference effects are observed. Any features, including defects and dirt, will generate a new series of interference pattern in the image. These interference “fringes” arise from the monochromatic nature and the coherence of the laser light. By rotating the mirror, these interference patterns also rotate within the image and due to averaging fringes are removed. Currently images are acquired 90 times within 3 seconds and averaged. As shown in Figure 2.9 (b), the resulting image quality has significantly improved.

The light reflected from the sample is transmitted through the commercial microscope (Reichert----Zetopan). A CCD camera is placed behind the objective to acquire the image. Software installed in our computer (LabView’s Vision development module) is capable of carrying out real time subtraction of a background from the image and subsequently displays them on the monitor. In practice, the sample is first magnetically saturated and the non-magnetic background recorded. We then perform the required action (such as changing field or introducing current pulse into sample) and
record image. Real time subtraction software enables magnetic signal to be directly monitored, which is usually not visible in the original un-processed image.

The real time subtraction suffers a major shortcoming: The background image obtained at the beginning of each scan is only good for 1 to 2 minutes. After that interference stripes, which are unrelated to the sample, appear on the subtracted image. The only way to remove these undesired stripes is to record a new background. This requires us to magnetically saturate the sample again. Therefore we only have an observation window which last for 1 to 2 minutes and only limited operations and number of images are possible to be taken within this time period. Complete automation of the whole experimental setup is thus required. The development of these interference stripes during the subtraction is likely to be related to the thermal instability of the laser. When temperature within the laser changes slightly, dimensions of the resonance cavity are slightly modified and lead to minor changes in the laser wavelength. Since the subtraction process is extremely sensitive to the laser wavelength any slight change will lead to incomplete subtraction and thus to the appearance of the stripes. Since the laser is most unstable immediately after it is started, development of interference stripes is most prominent in the first hour of operation of the laser. After a few hours, greater thermal stability is achieved and stripes develop much slower, enabling satisfactory Kerr measurements to be taken.
2.2.5 Comparison of MOKE with other magnetometry techniques

Generally, the sensitivity of MOKE is much higher than the VSM and SQUID “bulk” magnetometry techniques. Commercially available (e.g. Lakeshore) vibration sample magnetometer (VSM), features a noise floor of $10^{-7}$ emu at 10 seconds per data point. VSM collect signal form receiving coil placed in proximity to vibrating sample under a magnetic field and measure the total magnetic moment of the sample. The noise floor of the SQUID (Superconducting Quantum Interference Device) magnetometer manufactured by Quantum Design is $10^{-8}$ emu at ten seconds per data point. The Kerr hysteresis loop displayed in Figure 2.7 (d) is obtained on a permalloy sample with a light
spot of 50µm wide for 3 seconds per point. Assuming a SNR of 20:1 in this figure, we can calculate the noise floor of our MOKE magnetometer is about $10^{-9}$ emu, which is 10 times lower than the SQUID and 100 times lower than the VSM. In practice, Kerr imaging can even resolve magnetization moments down to areal dimension of $2 \times 2$ µm. This leads to a minimum resolvable moment to be below $3 \times 10^{-11}$ emu.

Magnetic Force Microscopy (MFM) is another powerful method to perform magnetic imaging. The resolution of Kerr microscopy (wide field) is ultimately limited by the wavelength of light (roughly half a micron). MFM typically can reach a spatial resolution down to tens of nanometer, an order of magnitude better than MOKE. However, MFM is sensitive only to the domain wall and not the local magnetization itself. Further, MFM relies on scanning a probe over a region, while Kerr microscope captures the profile of the whole area simultaneously. MFM is generally also slower than Kerr in achieving similar images.

**2.2.6 New trend in MOKE magnetometry and microscopy**

With the advent of femtosecond pulsed lasers, MOKE magnetometry has entered a “time-resolving” era [28]. “Pump and probe” method are being used. The “pump” light pulse is use to activate a photoconductive switch and to initiate an electrical current pulse that will high frequency magnetization oscillations in the sample. A “probe” light pulse (with controllable time delay) measures the magnetization vector using the Kerr effect. Successful attempts have been made to combine “time resolution” with Kerr microscopy in both wide field [29] and near field regime [30].
Other than promoting time resolution, improved spatial resolution is also valuable in Kerr microscopy. While the application of near field optical-scanning microscopy (NSOM) to obtain Kerr signal is mostly successful for the polar-effect [31], the longitudinal Kerr effect has also been recently investigated [32].

Figure 2.9: (a) Kerr image of 50×100µm channel without mirror rotation (b) With mirror rotation
3.1 Film deposition

Fabrication involves deposition of the film and their patterning using lithography methods. This section (3.1) is devoted to the discussion of film deposition, while the topic of lithography is discussed in section 3.2. In section 3.1, three types of deposition techniques that were used in this study are introduced. Since deposition is always on a supporting material----the “substrate”, we begin this section with substrate.

3.1.1 Substrate

Magnetic metallic microstructures are fabricated by “patterning” film deposited by various methods on a “substrate”. The “Substrate” is typically a bulk material that supports the thin film. It is often a high purity single crystal. The substrate surfaces are polished to “optical” quality such that their roughness is usually reduced to below 1 nm. Possible choices for substrates for this study are displayed in Table 3.1. Silicon wafer is popular for magnetic metallic film deposition. In applications that require electrical insulation from the substrate, the silicon surface is usually oxidized and generally
provides good adhesion of the magnetic metallic film while epitaxial growth with silicon is possible, polycrystalline film will result when deposition is on silicon oxide.

Magnesium oxide (MgO) is another choice of substrate for magnetic metallic films. They are insulating and have high thermal conductivity which is useful for high electrical current density applications where the high rate of heat dissipation in the substrate is desirable to reduce Joule heating. MgO is also transparent, which makes BLS and Kerr microscopy possible through the substrate on the the back side of the magnetic film possible. However, film adhesion to MgO is not generally as good as on silicon and proper caution must be taken during fabrication.

Zirconium oxide (ZrO$_2$) is another possible substrate choice although it has some inferior properties. First ZrO$_2$ is more brittle and makes subsequent patterning process more risky. ZrO$_2$ also has low thermal conductivity, making them undesirable for high current density applications.

Glass (conventional microscope slide) usually has excellent surface smoothness. However, the very low thermal conductivity also makes it a poor choice for high current density purposes.

Sapphire (Aluminum oxide or Al$_2$O$_3$) has properties very similar to MgO. Therefore sapphire is another popular choice for supporting magnetic metallic structures. However, sapphire suffers a major drawback: It is birefringent and therefore light becomes circularly polarized during propagation through the material. This inability to completely extinguish the resultant light by a polarizer makes Kerr microscopy from the backside of the film through a sapphire substrate impossible.
The insulating and hardness properties of diamond as well as extremely high thermal conductivity make it the best choice of substrate for high electrical current density applications. However, diamond substrates usually are manufactured by chemical vapor deposition (CVD) and have a relatively large density of defects. These defects act as light scattering centers and makes optical microscopy through diamond difficult. While high transparency is available in single crystal diamond, their high price make them impractical for wide usage in the laboratory.
<table>
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<th>Electrical conductivity</th>
<th>Optical Property</th>
<th>Hardness</th>
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<td>Brittle</td>
</tr>
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<td>Transparent, birefringent</td>
<td>Medium</td>
</tr>
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<td>Insulator</td>
<td>Transparent</td>
<td>Medium</td>
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<tr>
<td>Diamond</td>
<td>600 to 50000</td>
<td>Insulator</td>
<td>Transparent but usually hindered by scattering centers</td>
<td>Extremely hard</td>
</tr>
</tbody>
</table>

Table 3.1: Comparison of properties between different substrates suitable for magnetic metallic film deposition

### 3.1.2 Evaporation

As illustrated in Figure 3.1, the simplest method of film deposition is thermal evaporation. Solid source of the desired metal is usually placed in a tungsten boat. High density electrical current heats the boat until the temperature surpasses the melting point of the source material leading to its evaporation and subsequent deposition on the substrate. The time exposure of the substrate to the vapor determines the film thickness.
A vibration crystal detector calibrates the film thickness. It is obvious that the evaporation must be carried out under a vacuum, usually around 4 to $7 \times 10^{-6}$ Torr.

Thermal evaporation is relatively inexpensive and simple. It is good for metals with low melting points such as gold, copper, Aluminum and Chromium. Magnetic materials such as nickel (Ni) and iron (Fe) with high melting points are not suitable to be deposited by thermal evaporation as the tungsten boat cannot withstand the joule heating required to melt the material. Magnetic metallic structures are seldom deposited by thermal evaporation. Electron-beam evaporation is a more appropriate choice. Thermal evaporation is usually suitable for deposition of electrical contact pad for passing current. It is also suitable for depositing metal layer on top of resist before e-beam lithography for charge dissipation purpose (see section 3.3).

Figure 3.1: Schematic of a thermal evaporator
3.1.3 Diode sputtering

Ion beam diode sputtering (Figure 3.2) is another popular technique for depositing magnetic metallic films. While evaporation should be done in the highest possible vacuum, an argon gas of constant pressure ($8 \times 10^{-5} \text{Torr}$ in our case) is sustained during sputtering (after reaching a base pressure of $10^{-6} \text{Torr}$). When an ultrahigh potential difference (~1000 Volt) is introduced between the cathode and anode (distance between anode and cathode is approximately half a meter), electrons from the cathode move to the anode. During the transition, the electron may collide with and ionize an argon atom. The ejected electron will subsequently ionize another atom. Further, the electric field between the anode and cathode also helps to ionize the argon atom. All of these processes combine to form a chain reaction creating a plasma. $\text{Ar}^+$ ions in the plasma will be accelerated towards the cathode and bombard the target which is a bulk piece of the material to be deposited. Note that in Figure 3.2, the target is oriented at 45 degrees with respect to the accelerating direction of the $\text{Ar}^+$ ions. When the argon ion strikes the target surface, target atoms are ejected and subsequently reach the substrate. The $\text{Ar}^+$ ion regains its lost electron at the cathode to become neutral. It is the inertness of Argon which ensures that they do not react with the substrate. If a chemically reactive gas creates the plasma, undesirable mixed chemical composition of film will result.

The magnetic field created by electrical coils is utilized to focus the ion beam onto the target. However stray magnetic fields from these coils may introduce undesired magnetic anisotropy (preferred direction) on the film. Therefore the substrate is rotated during deposition to average out any magnetic field induced anisotropy.
The sputtering setup can be readily turned into an ion-milling machine by replacing the target with the sample to be milled. In this case the sample would face upright to the ion beam instead of the 45 degree tilt.

![Figure 3.2: Schematic of a diode sputtering machine.](image)

### 3.1.4 Magnetron sputtering

Magnetron sputtering represents an improvement over diode sputtering in providing an increased deposition rate. The main difference between magnetron and diode sputtering is that in the former the anode is placed beside the cathode, as depicted in Figure 3.3. Moreover, a permanent magnet placed beneath the target produces a magnetic field above, as represented by the dash line in the same figure. For the electron moving from cathode to anode, Lorentz force and cyclotron motion of the electron due to the magnetic field greatly increases its path length within the gas before reaching the anode. This resulting increased ionization yields a higher deposition rate.
In our case, a 300 Volt potential difference was applied across the anode and cathode (diameter of target and sputtering gun is several centimeters), with a current of 80 mA maintained between the two electrodes. A pressure of 4 milli-Torr of Argon gas was sustained within the deposition chamber. With the substrate approximately 15 to 20 cm above the target, 160 seconds was required to deposit 50 nm of Permalloy (Ni$_{80}$Fe$_{20}$) film on the substrate.

The situation depicted in Figure 3.3 is dc sputtering. In this case the target is required to be electrical conducting. For deposition of insulating materials rf (radio frequency) sputtering should be used instead. An ac voltage and current at 13.8 MHz is maintained between the two electrodes during the deposition. The use of ac current avoids charging the target which will stop further sputtering. Power consumption across the sputtering gun was maintained at 100 W and it took ~200 seconds to deposit 50 nm of nickel oxide (NiO). Rf sputtering can also be used to deposit conducting materials, such as Permalloy. However, in such a case dc sputtering is preferred due to the higher yield. In rf sputtering half of the time the Argon ions are moving away from the source and not sputtering the target.

Magnetic metallic film usually the deposited exhibits a certain anisotropy (directional preference of magnetization). This anisotropy is often determined by the surface profile of the substrate, crystal structure of the film (not applicable to polycrystalline film) and the magnetic field experienced by the film during deposition. If a magnetic anisotropy is desired, the substrate is placed in a magnetic field of several hundred Gauss with proper orientation.
3.1.5 Vacuum techniques

A vacuum environment is essential in the deposition process, since the source target atoms will require long a mean free path to reach the substrate. The extent of the vacuum also strongly determines the quality of the resulting film. Residual water vapor and oxygen inside the deposition chamber can easily contaminate the film and hamper the device performance. For instance, it is known that multilayers deposited in a vacuum of with pressure greater than $10^{-6}$ Torr do not display giant magnetoresistance effect. Therefore efforts are made to lower the base pressure.

3.1.5.1 Diffusion pump

The structure of a diffusion pump is illustrated in Figure 3.4. Oil at the base of the pump is heated and evaporates upward and forced to vent downward at several
locations on the upper section of the pump. The oil vapors collide with air molecules and force them downward. The air molecules reaching the bottom of the pump are subsequently removed by the mechanical pump. The diffusion pump is always backed up by a mechanical pump and is never exposed to atmospheric pressure to avoid the oil boil and burn. The oil vapor inside the pump will eventually condense at the bottom and restart the cycle.

The liquid nitrogen trap is condenses any residual water vapor inside the deposition chamber to further lower the pressure. The trap also prevents any oil vapors in the diffusion pump from entering and contaminating the deposition chamber.

The mechanical pump alone can attain a pressure down to 50 milli-Torr. With the diffusion pump and the cold trap, a pressure at the level of $10^{-6}$ Torr can be achieved.

![Diffusion pump diagram](image)

Figure 3.4: Diffusion pump
3.1.5.2 Turbomolecular pump

The turbomolecular pump is one of the most popular approaches to attain ultra-high vacuum (UHV). It is clean (oil free) and reliable. Figure 3.5 displays the blades of a turbomolecular pump turbine. Given that the pressure is below certain level, the blade spacing is designed to be smaller than the mean free path of air molecules. In such a circumstance, intermolecular collisions are rare and the molecules essentially collide only with the pump blade. With a rotation speed of 30 krpm (kilo rotation per minute), a vacuum of $10^{-10}$ Torr can be achieved. However, the performance of turbomolecular pump are very bad (and dangerous) at atmospheric pressure and unsuitable for chamber pressures above $10^{-4}$ Torr.

![Figure 3.5: Picture of blades of a turbomolecular pump. (Adopted from Ref. 33)](image-url)
3.1.5.3 Sublimation pump

A sublimation pump is another approach to improve the vacuum within the chamber. The vapor from heated Titanium inside the chamber will condense (without going through the liquid phase) directly into titanium solid on the inner-wall of the chamber which gives the name “sublimation”. Since Titanium is a reactive metal it will chemically react and absorb residual atmospheric molecules within the chamber during sublimation.

3.1.5.4 Baking of deposition chamber

The inner wall of the deposition chamber is constantly out-gassing slowly releasing molecules such as oxygen and water, and hinders the chamber reaching the required vacuum. By switching on a high Intensity lamp inside the chamber, the absorbed H$_2$O and O$_2$ molecules are “baked” off and released in a short period of time. Once released these molecules are pumped out of the chamber by the turbomolecular pump.

3.2 Photolithography

Lithography is the process of transferring of a pre-designed pattern to a miniature scale. Except for self-assembled (bottom-up) system (life is a very good example of a self-assembled system!), lithography is the only main way to achieve human designed and control at the micro and nano-scale. Miniaturization is achieved by the projection of an energy carrying beam onto a targeted sample. Different beam of projection lead to lithography prefixed with different names, such as: photo-, e-beam-, focused ion-beam-lithography etc. In this section, we concentrate on photolithography.
3.2.1 Subtractive techniques

Lithography usually transfers a pattern on to sample by means of subtractive or additive strategy. Figure 3.6 illustrates a series of steps involved in a subtractive mode. The technique is called subtractive since the film is first deposited on the substrate before the pattern is formed by etching away undesired regions (subtraction of the film). If the pattern is formed by lithography resist before the film is deposited, then the technique is named “additive”. In Fig. 3.6 (a), a single or multilayer magnetic metallic film is deposited on to the substrate. A layer of “photoresist” is then spun on to the top of the sample (Fig. 3.6 b). The photoresist is a fluid of polymers, usually red in color. The word “photo” refers to their sensitivity to light and “resist” describes its ability to resist corrosion against strong acid. Droplets of photoresist are placed onto the sample surface and the sample is rotated at a speed of several thousand rpm. The spinning evenly distributes several micrometer thickness of the photoresist. The photoresist used in this study was “AZ 4620” which is diluted with “thinner” solution to make the resulting photoresist fluid less dense and thus eventually achieving a thinner resist layer.

Once the resist layer is deposited, the sample is baked prior to the step of exposure. Fig. 3.6 (c) schematically illustrates appropriate locations of the resist on the surface that are exposed to light. Polymers of the resist exposed to light disintegrate into smaller molecules. After exposure, the sample with the exposed resist is immersed into the developer solution (400K) (Fig. 3.6 d). Regions previously exposed to light preferentially dissolve in the developer solution, while the unexposed area still composed of large polymers remain insoluble. The required development time is usually 3 to 4
minutes. After the exposed resist is completely removed in the development process, the sample is subject to “etching”. There are basically two kinds of etching: “dry etch” and “wet etch”. In a wet etch process, the sample with the resist is immersed into a strong acid solution. The acid etches the sections of the film exposed to the acid, while the film under the photoresist remains protected ---- Fig. 3.6 (e). Finally, the remaining resist is removed by a solvent such as acetone (Fig. 3.6 f). The overall effect is the transferal of the desired pattern onto the magnetic metallic film.

In this study, the following formula in the wet etch process was found to be exceptionally useful: $200\text{ml } \text{H}_2\text{O} + 25\text{ml } \text{H}_2\text{SO}_4 (95.9\%) + 7\text{ml } \text{H}_2\text{O}_2 (30\%) + 3\text{ml } \text{HF (49\%)}$. The etch rate on Permalloy is roughly 10 nm per second.

The wet etch process always has the drawback of possibly undercutting the resist, i.e. etching into film area covered by the photoresist. Also uncontrollable randomness enters the etching process leading to rough boundaries. The roughness usually is of the order of a micron. Wet etch, although fast and easy, basically is not applicable in achieving features with dimensions smaller than one micron in the case of metallic magnetic microstructure fabrication.

Ion beam milling is a good example of dry etching. The technique of diode sputtering (section 3.1.3) can be easily converted into an ion beam mill. Care has to be taken to ensure the film in the exposed region be milled off completely before the photoresist is fully removed, otherwise the film in region protected by the resist will be etched, which is undesired. Typically, with an accelerating voltage of 1 keV, argon ions remove tens of nanometer of resist per minute. Etch rates of common metals under the same condition are roughly similar. However, during milling the resist is not
only being etched vertically, but also horizontally. Therefore a ten micron wide channel can finally end up as 8 micron. It was found that the resist can better protect the film if the milling is done at cryogenic temperatures. A cold stage should be used to keep the substrate at low temperature to avoid any excessive heating associated with milling.

![Figure 3.6: Steps of photolithography with subtractive techniques](image)

Figure 3.6: Steps of photolithography with subtractive techniques (a) Deposition (b) Spinning of photoresist (c) Exposure (d) Development (e) Etching and (f) Stripping off photoresist.

3.2.2 Contact versus projection

There are two types of exposure: contact and projection. As depicted in Figure 3.7 (a), in contact exposure, the mask, which is fabricated on a piece of glass, is brought into contact with the photoresist. The mask is usually made by e-beam lithography. The light source for photo exposure is placed on top and Figure 3.7 (b) illustrates a projection lithography setup. The mask is held behind the lens and with the help of a light source, a diminished shadow image of the mask is produced on the resist surface.
Most advanced commercial optical microscopes have the ability to perform projection exposure. An Olympus Optical Microscope (model BH-2) (displayed in Figure 3.7 c) was used to achieve a reduction ratio of 6:1 in this thesis project. The arrow in Fig. 3.7 (c) indicates the location to insert a mask for projection lithography. The mask was made with a second round of projection with reduction ratio of 20:1 giving rise to an overall reduction of 120:1. Thus given that the minimal features from a commercial laser printer are 1 mm wide, the lithography routine can achieve structures on the sample down to 8 micron.

Contact lithography is preferred in conventional chip manufacturing since the mask can cover a wide area (the whole silicon wafer), while the working area of projection lithography is largely limited by the field of view of the microscope objective. In the case of the highest magnification, the field of view can be decreased down to hundreds of microns. However, projection lithography is preferred in a research process since it does not require contact of a rigid mask with the sample. Such contact is often risky and may damage delicate structures on the sample surface. Projection lithography is empirically more stable than contact lithography. In contact lithography, multi-reflection of light between the sample surface and the mask at the pattern edges often lead to blurring of the eventual pattern boundary. Special care must be taken to deal with this problem. The ultimate resolution limit of photolithography is limited by the wavelength of light, which is in the submicron region. Electron beam lithography is preferred in managing tasks at this level.
3.2.3 Additive techniques

The subtractive method (section 3.2.1) suffers a major disadvantage: The etching step usually leads to the undesired side effect wherein the acid used during wet-etch contaminates the sample. In dry etching, overshooting is required to ensure the film is completely removed. It is therefore necessary to ion mill beyond what is required for considerable depth into the substrate. This destroys the substrate surface for further fabrication treatment. Therefore although in early years subtractive method used to produce patterned magnetic metallic films [34], the present trend is towards the additive approach [35].

Procedures in additive photolithography are displayed in Figure 3.8. A layer of photoresist is first spun on top of the bare substrate (Fig. 3.8a). Baking, exposure and
development, as described in section 3.2.1, is carried out. This leads to the situation depicted in Fig. 3.8(b). In (c), a magnetic metallic film is deposited on top of the “substrate plus resist” surface. In (d), the entire sample is immersed in an acetone solution which dissolves any remaining photoresist while removing the magnetic film on top at the same time thereby leaving the desired magnetic metallic film micro-structure on the surface.

Compared to the subtraction technique, the additive method is cleaner and the substrate surface quality is retained. However, caution must be taken in the lift-off step. Since, the deposited metallic films on the sides with and without photoresist are connected to each other (see Fig. 3.8e), when the side of the film residing on the resist is removed, it could also easily tear the other side away from the sample. The lift off procedure should thus be done slowly, with the acetone solution being stirred only moderately. Some articles suggest the use of ultrasound to remove the resist. Our experience however suggests this is not a good practice since adhesion of the film may not be strong enough to withstand the vibrations associated with the ultrasound.

After the lift off procedure, it is possible that a thin layer of resist remains on the sample surface. This layer cannot be stripped away by solvents. This is especially true for resists that have been baked over long time periods. In such circumstances, the resist should be removed by ozone cleaner instead. Weak ozone cleaners are able to selectively remove organic matter (the resist) while leaving the substrate and the magnetic metallic film undamaged.
3.3 Electron beam lithography

The ultimate limitation of photolithography is the wavelength of the light used to expose the resist. Due to diffraction features smaller than the wavelength cannot generally be resolved. In electron beam (e-beam) lithography, the much smaller de Broglie wavelength of the fast moving electron results in a much higher resolution than photolithography, with features tens of nanometers being feasible.
Figure 3.9 depicts an e-beam writing process on the top of “electron beam resist”. The resist for e-beam lithography is PMMA (Polymethyl methacrylate). During the writing process, the movement of the beam on the sample is controlled by a computer software (J. C. Nabity Lithography System), which is pre-programmed to write a desired pattern. The beam motion is actuated by deflecting the e-beam with an applied electric field. The precise control of the electron beam position on the sample surface in e-beam lithography can be realized from the fact that the electrons are moving at high speed, and the very large applied electric field gives rise to only nanometer scale electron deflection. The zig zag path in Fig. 3.9 represents the path of the electron beam, which results in the “cross” pattern as shown. The pattern will not develop properly unless the correct electron dosage is utilized. Dosage is the quantity of charge received by the PMMA per unit area, determined by the magnitude of current associated with the electron beam, the speed of movement of the e-beam and the distance between adjacent tracks of the zig zag path. The required dosage also depends on the substrate, as a large fraction of bond breakage in the PMMA is due to electrons rebounding from the substrate. Thus different electron reflectivities of different substrates will influence the optimum dosage. Further, the dosage will also depend on the details of the structural features. Patterns with narrow widths often require higher dosage. Therefore, before writing a large pattern, it is recommend to first test a small representative segment of the pattern to determine the optimum dosage.

PMMA with molecular weight equals 496K was the resist of choice in this thesis. By spinning the sample at 2500 rpm, a PMMA thickness of roughly 300 nm was obtained. The PMMA coated sample was baked at 180°C for one hour. The required e-beam
dosage was usually around several hundred µC/cm² at an acceleration voltage ~ 20 to 30 kilo-volt. The written sample was then developed in 1:3 MIBK:IPA (isopropyl alcohol) solution for 20 seconds. As noted, the smallest achievable feature depends on the smallest possible e-beam spot size on the sample. This in turn is highly dependent on the quality of the SEM in use. For the electron microscope used in this study (JEOL JSM-5400), the minimal feature size achievable was half a micron. A higher quality SEM usually has a higher acceleration voltage (up to 100 kV), that enables achieving features of smaller dimension.

Care must be taken when the substrate is not electrically conductive, since charging of the surface can occur during the writing process. The charging can in turn deflect the electron beam path and lead to distortion of the resultant pattern. A 10 nm thick aluminum layer should be deposited on the top of the PMMA before the writing process to avoid charging. The particular choice of metal to deposit for the purpose of charge dissipation is important. It is known for example that if gold is used instead of aluminum, failure will result due to intermixing of gold with the PMMA and formation of a hard intermediate layer. Such a hard layer leads to problems during development. Figure 3.10 displays an AFM (atomic force microscope) image of a 2 micron wide, 50 nm thick permalloy wire fabricated on a SiO₂/Si surface, using the e-beam lithography with the aforementioned additive method.
Figure 3.9: Schematic illustration of electron beam lithography

Figure 3.10: AFM image of a 2 micron wide, 50 nm thick permalloy wire fabricated on the surface of SiO$_2$/Si.

54
CHAPTER 4

TUNABLE FIELD IMBALANCE AND SPIN PRECESSION

4.1 Introduction

Thin ferromagnetic/normal metal multilayer structures have received much attention over the last two decades and have emerged as magnetic field sensors, computer hard-disk read heads, and magnetic random access memory elements [36]. Although initially the current density for these applications was retained at modest levels, it has been raised to values beyond $10^7$ A/cm$^2$ in order to enhance signals [37], produce controlled movement of domain walls [38], and create measurable spin–torque transfer between magnetic layers [39]. In these studies, attention was largely focused on the influence of the current on the static component of the magnetization for memory applications. As witnessed by the recent observation of microwave emission from structures in a current perpendicular to plane geometry [40], there is growing interest in exploring the dynamic magnetization component. The microwave emission in these systems has led to suggestions that it could be utilized in mesoscopic oscillators and electronic clocks. Such technologies would require a more detailed understanding of the interactions between the dynamic components of the magnetization and the required high current densities. While these investigations have generally been limited
to the spin–torque transfer mechanism, the impact of the Amperian self-magnetic field
\( (H_{\text{cur}}) \) associated with the current on the precessing magnetization has received little
attention. Interestingly, for a current-in-plane geometry, an external in-plane field \( (H_{\text{ext}}) \)
together with such Amperian fields at current density values of \( 10^7 \, \text{A/cm}^2 \) fall in a range
that generates a spatial field imbalance within a multilayer structure to dramatically
modify the oscillation profile of the magnetization. In this section, we present a Brillouin
light scattering (BLS) study providing direct measure of such local variations of the
dynamic magnetization components in distinct spatially separated layers. The results,
while in agreement with first-principle calculations, are also analyzed within a
mechanical analog of a coupled pendulum in an unbalanced gravitational field.

4.2 Methodology

BLS experiments were performed on metallic trilayer strips carrying electric
current. A permalloy-copper-permalloy trilayer of \( \text{Ni}_{80}\text{Fe}_{20} \) (\( d=30 \, \text{nm} \))/Cu (\( d_o=45 \, \text{nm} \))/\( \text{Ni}_{80}\text{Fe}_{20} \) (\( d=30 \, \text{nm} \)) was sputter deposited on a diamond substrate (Sumitomo
Electric Industries, Japan), followed by photolithography and wet etching to yield the
patterned channel (width and length 50 and 100 mm, respectively) illustrated in Figure
4.1. The sample is named “sample #1”. The selection of diamond was based on its
superior thermal conductivity to minimize effects associated with resistive heating. The
Cu spacer thickness (\( d_o \)) of 45 nm enables passage of a relatively large current through it,
while providing the framework for the spin waves in the two NiFe layers to be coupled.
The current, \( I_{\text{channel}} \), was introduced to the narrow channel through leads placed outside
the constriction leading into the patterned segment and current densities as high as
56
$4 \times 10^7$ A/cm$^2$ were realized. The electrical conductivity of Cu deposited on separate substrates under identical conditions was measured to be 2.2 $\mu\Omega \text{cm}$; that of permalloy was 20 $\mu\Omega \text{cm}$, resulting in 88% of the current to flow through the Cu spacer for the specific structures utilized in this study. For simplicity of the analysis described below, we neglect effects of the small current flow outside of the Cu layer. An in-plane external field $H_{\text{ext}}$ was applied as shown in Fig. 4.1, perpendicular to the direction of current flow. Brillouin spectra ($p\rightarrow s$ scattering) were recorded as a function of $I_{\text{channel}}$ in backscattering utilizing 70 mW of 514.5 nm radiation at a low temperature of 25 K (to minimize effects due to Joule heating). The laser probe beam was incident at 45° to the surface in a plane perpendicular to the external field with the beam focused to 50 $\mu$m diameter, and typical accumulation times of 20 min.

![Figure 4.1: Schematic sketch of experimental setup. The arrow on the circuit defines the direction of positive current. Inset: Cross section of magnetic double layer with coordinate axes. Current is along the $+y$ direction.](image)

57
4.3 Preliminary results

We present results for two limits of the external applied magnetic field; the high \( (H_{\text{ext}} = 600 \text{ Oe}) \) and low \( (H_{\text{ext}} = 50 \text{ Oe}) \) field cases. Fig. 4.2 illustrates Brillouin spectra recorded as a function of the channel current \( (I_{\text{channel}}) \) when \( H_{\text{ext}} \) is fixed at 600 Oe that results in parallel moments and a saturated structure. For zero current three Stokes modes are observed; peak (V) at 18.6GHz is identified as the volume mode while those at 12.0GHz and 14.5GHz are the surface optic (SO) and surface acoustic (SA) modes. It is noted that the volume mode depends strongly upon the intralayer exchange while, as discussed below, the surface modes are satisfactorily described in the dipolar approximation. These mode assignments are based on their observed frequency dependence on the angle of incidence (not shown) and agreement with calculations described below. With the introduction of a channel current of +1.0A (Fig. 4.2b), two primary changes from the \( I_{\text{channel}} = 0 \text{A} \) spectra are evident. These are: (a) reduction of the frequency of the volume mode from 18.6 GHz to 18.0 GHz and (b) an enhancement (suppression) in the SO (SA) peak intensity, resulting in changes to the relative intensities of the in- and out-of-phase coupled modes SA and SO; the frequencies of the pair of surface excitations are largely unaffected by \( I_{\text{channel}} \) for \( H_{\text{ext}} = 600 \text{Oe} \). Upon reversing the current (from \( I_{\text{channel}} = +1.0 \text{ A} \) to \(-1.0 \text{A}\)), the volume mode is observed to harden reverting to 18.6 GHz, its zero current value (Fig. 4.2c). The current switch also leads to suppression (enhancement) of the SO (SA) peak intensity. The reversibility of these changes in the Brillouin peaks is confirmed in Figs. 4.2d, e and f when \( I_{\text{channel}} \) is once again switched prior to being reduced to zero.
Figure 4.2: Brillouin spectra recorded at fixed external field ($H_{\text{ext}}$) of 600Oe as a function of current. The dash line is a guide to the eye.

Explanations for these spectral changes with $I_{\text{channel}}$ are traced to the accompanying Amperian fields ($H_{\text{cur}}$) that orient in opposite directions in the top and bottom permalloy layers. The largest field generated in our study is estimated to be ~100Oe for $I_{\text{channel}}$~1 A. Since the surface spin wave extends over both NiFe layers, the oppositely directed $H_{\text{cur}}$ fields in each magnetic layer have, as observed, little overall influence on the surface magnon frequencies (for $H_{\text{ext}}$=600Oe). While Joule heating associated with $I_{\text{channel}}$ will lead to a decrease in the saturation magnetization, we estimate
such thermal contributions to be minimal accounting, at most, for a \( \sim 0.3 \)GHz drop in the
frequency \((f_{SO} \text{ and } f_{SA})\) of modes SO and SA. These changes lie within the error margins
of the measurements. On the other hand the frequency of the volume mode that is
confined to each NiFe layer, reveals a finite response to \( I_{\text{channel}} = +1.0 \) A in the presence
of \( H_{\text{ext}} = 600 \) Oe. In this case the magnon softens for positive current since the effects
due to \( H_{\text{cur}} \) and Joule-heating augment one another in the top NiFe layer where the probe
laser beam is incident. This leads to an overall 0.6GHz suppression of the mode (Figs.
4.2b and d). For reversed current flow \( (I_{\text{channel}} = -1.0 \) A), the \( H_{\text{cur}} \) driven frequency
increase in the upper NiFe layer is counter balanced by effects of heating that leave the
volume mode frequency essentially unchanged (Figs. 4.2c and e). This interpretation of
compensating effects between Amperian field and Joule heating is also consistent with
results from separate measurements, observing (not shown) independently the field and
temperature dependence of the volume mode frequency.
Figure 4.3: Schematic illustration of in-phase (SA) and out-of-phase (SO) magnetization oscillations shown on precession cones for various Amperian ($H_{cur}$) and external field ($H_{ext}$) directions and magnitude in case of $|H_{ext}| > |H_{cur}|$ (parallel alignment of magnetization). The rectangles represent the two magnetic layers in cross-section. The probe laser beam is incident from the left and directions of $I_{channel}$ indicated by ⊗ and ⊙.
To account for the variations of the Brillouin intensity with $I_{\text{channel}}$, we trace into a simplified version of change in SA and SO spin wave amplitude profile in the presence of the spacer current, as illustrated in Fig. 4.3. More rigorous analysis based on solving Maxwell and Landau-Lifshitz equations is presented later in this section. Fig. 4.3a and b shows the oscillation of SA and SO spin wave in case of $I_{\text{channel}}=0$ respectively. In this simplified version of the mode profile, the oscillation amplitudes for the two spin wave modes on the left magnetic layer, where the light is scattering, are identical. This corresponds to the spectrum of Fig. 4.2a. More rigorous calculations account for the unequal peak height of SA to SO mode in Fig. 4.2a. We ignore the inequality at present.

In Figure 4.3c and d, current flows out of paper through the spacer layer and leads to enhancement (suppression) of the SO (SA) amplitude on the left magnetic layer, thus accounting for the enhancement (suppression) of SO (SA) peak intensity in Fig. 4.2b and d. The enhancement (suppression) of the SO (SA) amplitude on the left magnetic layer is associated with modification of the spin wave mode profile across the bilayers (Fig. 4.3c and d). This is explained rigorously later in this section and more intuitively by a coupled pendulum model in section 4.5. In Figure 4.3e and f, current flows into the paper through the spacer layer and leads to suppression (enhancement) of the SO (SA) amplitude on the left magnetic layer, thus explaining the suppression (enhancement) of the SO (SA) peak intensity in Fig. 4.2c and e.

To rigorously account from first principle the changes of SA and SO spin wave amplitude profile with respect to the spacer current, we solve for the spin wave modes of a ferromagnetic double layer in the presence of an in-plane spacer layer current, in a
fashion similar to the zero current case [41]. The inset to Fig. 4.1 illustrates the coordinate axes with the external applied field $H_{\text{ext}}$ directed along $+z$. The free current density is given by

$$J_f = \begin{cases} J_\circ \hat{y}, & \text{for } |x| < 0.5d_o \\ 0, & \text{otherwise} \end{cases}$$

where $J_\circ$ is a constant (independent of spatial position). The net magnetic field satisfying

$$\nabla \times \mathbf{H} = \frac{4\pi J_f}{c}$$

is given by $\mathbf{H} = H_{\text{eff}}(x)\hat{z} + \mathbf{h}(x,y)e^{i\omega t}$, where

$$H_{\text{eff}}(x) = \begin{cases} \frac{4\pi J_\circ x}{c} + H_{\text{ext}}, & \text{for } |x| < 0.5d_o \\ -H_{\text{cur}} + H_{\text{ext}}, & \text{for } x > 0.5d_o \\ H_{\text{cur}} + H_{\text{ext}}, & \text{for } x < -0.5d_o \end{cases}$$

such that

$$H_{\text{cur}} = \frac{2\pi J_\circ d_o}{c}$$

and $\mathbf{h}(x,y)$ is the fluctuating component of $\mathbf{H}$. The magnetization is similarly given by

$$\mathbf{M} = \begin{cases} M_\circ \hat{z} + \mathbf{m}(x,y)e^{i\omega t}, & \text{for } 0.5d_o < |x| < 0.5d_o + d \\ 0, & \text{otherwise} \end{cases}$$

Introducing a magnetic potential $\psi(x,y) = X(x)e^{iky}$ where $\mathbf{h} = \nabla \psi$, it follows that the Landau-Lifshitz equation $d\mathbf{M}/dt = \gamma(\mathbf{M} \times \mathbf{H})$ yields
\[ 4\pi m_x = \kappa \frac{\partial \psi}{\partial x} - iv \frac{\partial \psi}{\partial y}, \]
\[ 4\pi m_y = iv \frac{\partial \psi}{\partial x} + \kappa \frac{\partial \psi}{\partial y}, \]  

where

\[ \kappa = \frac{\Omega_H}{\Omega_H^2 - \Omega^2}, \quad \nu = \frac{\Omega}{\Omega_H^2 - \Omega^2}, \quad \Omega_H = \frac{H_{\text{eff}}(x)}{4\pi M_o}, \quad \Omega = \frac{\omega}{4\pi\gamma M_o} \]  

The parameters \( \kappa \) and \( \nu \) hence have different values in the two magnetic layers. Writing \( X(x) \) as

\[ X(x) = \begin{cases} 
  e^{kx}, & \text{for } x < -(0.5d_o + d) \\
  B_1 e^{kx} + B_2 e^{-kx}, & \text{for } -(0.5d_o + d) < x < -0.5d_o \\
  C_1 e^{kx} + C_2 e^{-kx}, & \text{for } -0.5d_o < x < 0.5d_o \\
  D_1 e^{kx} + D_2 e^{-kx}, & \text{for } (0.5d_o + d) < x < 0.5d_o \\
  E_1 e^{kx} + E_2 e^{-kx}, & \text{for } x > (0.5d_o + d) 
\end{cases} \]  

and enforcing the boundary conditions associated with the continuity of \( \psi(x,y) \) and the normal component of \( B \) at the interfaces result in:

\[ (1 + \kappa) \frac{dX^i}{dx} + \kappa v X^i = \frac{dX^e}{dx} \]  

\[ X^i = X^e \]  

where superscripts \( i \) and \( e \) refer to the interior and exterior of the magnetic layers. For the known parameters \( d(30 \text{ nm}), d_o(45 \text{ nm}), H_{\text{exx}}(600 \text{ Oe}), H_{\text{cu}}(0 \text{ or } \pm 100 \text{ Oe}), \) wavevector \( k (+1.73 \times 10^7 \text{ m}^{-1}) \) and \( M_o(812 \text{ Oe}), \) definite numerical values for the coefficients \( B_1, B_2, C_1, C_2, D_1, D_2, E_1 \) and \( E_2 \) can be deduced from equation 4.5 applied at each magnetic-non-magnetic material interface for a given \( \omega \). Further, requiring the coefficient \( E_1 \) to vanish
for a bounded solution of $\psi$ leads to two solutions of $\omega$ corresponding to modes SA and SO \[42\]. Their mode profiles $(m_x, m_y)$ as a function of $x$ are then determined from equation 4.2. The results are illustrated in Fig. 4.4a, b and c, corresponding to the three cases: zero current, current $J_f$ along $\pm y$ direction (leading to a current induced field of $H_{cur} = \pm 100$ Oe). In Fig. 4.4, the spin wave is propagating in the $-y$ direction (out of page) and the external field ($H_{ext}$) is along $+z$, while the $x$ direction is perpendicular to the bilayer surface. It is evident from Fig. 4.4 that the SO spin wave mode amplitude (solid line) within the NiFe layer located at negative $x$ is, in comparison to the zero current case of $\{m_x, m_y\} = \{0.92, -1.72\}$, enhanced to $\{m_x, m_y\} = \{1.08, -1.92\}$ (suppressed to $\{m_x, m_y\} = \{0.77, -1.52\}$) when the current flow in the same (opposite) direction as the spin wave propagation. On the other hand that the SA mode amplitude at negative $x$ (dashed line in Fig. 4.4), in comparison to the zero current case of $\{m_x, m_y\} = \{1.28, -1.57\}$, is instead suppressed to $\{m_x, m_y\} = \{1.17, -1.39\}$ (enhanced to $\{m_x, m_y\} = \{1.36, -1.71\}$) by the same corresponding current flow thus accounting for changes in Fig 4.2 in the relative Brillouin intensities of the two surface modes SA, SO with $I_{channel}$. In contrast to the peak intensities, the calculated frequencies $f_{SO}$ and $f_{SA}$ of the surface waves in Fig. 4.4 are largely unaffected by the current. This behavior hence explains the insensitiveness of surface wave frequencies to $I_{channel}$ (Fig. 4.2).
Figure 4.4: Calculated mode profiles of SO mode (solid lines) and SA mode (dotted lines) for different free current density $J_f$ flowing in the Cu spacer layer. The mode frequencies $f_{SO}$ and $f_{SA}$ of the two surface spin waves are also indicated for each case. The external field ($H_{ext} = 600\text{Oe}$), is much larger than the current induced field $H_{cur} \sim 100\text{Oe}$. The direction of spin wave (SW) (k vector) is out of the page (-y) and the external field ($H_{ext}$) is along +z. All mode profiles have been normalized to include the same spin wave energy (integration of $|m_x|^2 + |m_y|^2$ over $x$).

The results discussed above were recorded for external field of $H_{ext} = 600\text{ Oe}$ that far exceeds the maximum current induced field ($H_{cur}$) of $\sim 100\text{ Oe}$. Measurements recorded with $H_{ext}$ at nominally 50 Oe - where the current induced field now exceeds the external field - are illustrated in Fig. 4.5. In the absence of $I_{channel}$, the three spin
waves V, SA and SO appear as symmetric Stokes and anti-Stokes excitations, a behavior consistent with the parallel alignment of the soft NiFe layer due to the modest external applied field (Fig. 4.5a). In the presence of $I_{\text{channel}} = +1.0$ A, the Stokes and anti-Stokes frequencies of the surface modes become strongly asymmetric (Fig. 4.5b) - a behavior reflecting anti-parallel alignment of the magnetizations of the two permalloy layers [43]. This frequency asymmetry between Stokes and Anti-Stokes surface spinwaves is associated with different frequency splitting for waves traveling in opposite direction due to different coupling strength, which will be accounted later in this section. Concurrently, the intensity of the anti-Stokes SA peak is stronger than the corresponding anti-Stokes SO mode. Upon reversing $I_{\text{channel}}$ from +1.0A (Fig. 4.5b) to −1.0A (Fig. 4.5c), the Stokes and anti-Stokes spectra reverse. This reversal is in agreement with the NiFe magnetization continuing to be aligned anti-parallel, wherein the magnetization in each layer was separately reoriented when current flow is reversed. In this case the relative strength of the SO and SA peaks also switch subsequent to the current reversal, where the Stokes peak SO is now stronger than the Stokes SA peak (Fig. 4.5c). Upon reducing $I_{\text{channel}}$ to zero, the magnetizations are now influenced by $H_{\text{ext}} = 50$ Oe that leads to saturation and parallel alignment consistent with the symmetric Stokes, anti-Stokes frequencies illustrated in Fig. 4.5d.
Figure 4.5: Brillouin spectra recorded with a fixed external field $H_{\text{ext}} = 50$ Oe as a function of $I_{\text{channel}}$. 
Figure 4.6: Schematic illustration of in-phase (SA) and out-of-phase (SO) magnetization oscillations shown respectively on precession cones for various Amperian ($H_{cur}$) and external field ($H_{ext}$) directions and magnitude in case of $|H_{ext}| < |H_{cur}|$ (anti-parallel alignment of magnetization). The rectangles represent the two magnetic layers in cross-section. The probe laser beam is incident from the left and directions of $I_{\text{channel}}$ indicated by $\otimes$ and $\bigcirc$. 

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<td>SO mode</td>
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Changes to spectra displayed in Fig. 4.5 are traced from the simplified version of modification of SA and SO spin wave precession with respect to different external and current induced field as illustrated in Fig. 4.6. The spectrum of Fig. 4.5a with zero $I_{\text{channel}}$ corresponds to the case of Fig. 4.3a and b, with parallel alignment of magnetization on the two layers. Fig. 4.5b where $I_{\text{channel}}=+1.0A$ corresponds to the case of Fig. 4.6a and b, with a spacer current flowing out of the paper. The Amperian field associated with the current reverse the magnetization on left layer (where light is scattering) leading to anti-parallel alignment of magnetization. Since the amplitude of SA mode is larger than SO on the left layer (Fig. 4.6a and b), the SA peak intensity is higher than SO in the spectrum of Fig. 4.5b. In the spectrum of Fig. 4.5c, $I_{\text{channel}}$ is switched to -1.0A. This corresponds to the case of Fig. 4.6c and d with a spacer current flowing into the paper. The magnetizations of the two layers are both reversed such that anti-parallel alignment is retained. Since the amplitude of SO mode is now larger than the SA on the left layer (Fig. 4.6c and d), SO peak intensity is higher than SA in the spectrum of Fig. 4.5c.

Result from first principle calculations based on Maxwell and Landau-Lifshitz equation are as follow: For the situation $H_{\text{cur}} > H_{\text{ext}} = 50$ Oe (Fig. 4.5), the spin wave behavior and the orientations of $M_o$ are now determined primarily by $H_{\text{cur}}$. The results for the spin wave profiles are displayed in Fig. 4.7 with calculated SO, SA frequencies in good agreement with the experimental asymmetric Stokes, anti-Stokes surface wave frequencies (Fig. 4.5b and 4.5c). This signature of anti-parallel alignment of magnetization of the two magnetic layers is illustrated by different values of $f_{\text{SO}}$ and $f_{\text{SA}}$ in Fig. 4.7b and 4.7c. This asymmetry in the surface spinwave frequencies is associated with different effective magnetostatic interaction between oscillations on the two
magnetic layers due to the different mode profile. In Fig. 4.7b, the oscillation amplitudes of spinwave mode are larger on the inner surfaces of the two magnetic bilayers (the interface between the magnetic layers and the central copper spacer) when compared to the outer surfaces (interface between the magnetic layers and air). This lead to stronger mode coupling strength between the layers and therefore larger frequency splitting ($f_{SO} - f_{SA} = 17.4 \text{ GHz} - 7.3 \text{ GHz} = 10.1 \text{ GHz}$). In the opposite case, in Fig. 4.7c, the oscillation amplitudes of the spinwave modes are smaller on the inner surface of the two magnetic bilayers than on the outer surfaces. This lead to weaker mode coupling strength between the layers and therefore smaller frequency splitting ($f_{SO} - f_{SA} = 12.9 \text{ GHz} - 11.8 \text{ GHz} = 1.1 \text{ GHz}$). The large difference in frequency splittings (10.1GHz versus 1.1GHz) leads to the Stoke and Anti-Stoke surface spinwave frequency asymmetry in Fig. 4.5b and c.

Figure 4.7 also provides an explanation for the reversal of SO and SA mode intensities in the switch from the anti-Stokes spectra of Fig. 4.5b to Stokes in Fig. 4.5c. This switch occurs since in Fig. 4.7c, within the NiFe layer located at positive x the amplitude of the SA mode ($\{m_x, m_y\} = \{1.37, 2.18\}$) is stronger than that of the SO mode ($\{m_x, m_y\} = \{-0.98, -1.33\}$), while the reversed behavior holds true in the NiFe layer at negative x (with SO amplitude $\{m_x, m_y\} = \{1.36, -1.92\}$ versus SA amplitude $\{m_x, m_y\} = \{0.97, -1.61\}$).
Figure 4.7: Same as Fig. 4.4 except the external field $H_{ext} = 50$ Oe, that is smaller than the current induced field $H_{cur} = 100$ Oe. The direction of the overall magnetization $M_o$ is thus primarily determined by $H_{cur}$.

### 4.4 Field dependence results

Figs. 4.8a-f show representative anti-Stokes Brillouin spectra for different $H_{ext}$ and $H_{cur}$ (derived from $I_{channel}$ and channel width) taken from another sample, with a channel dimension of width and length $70 \times 150 \mu$m (sample #2). Three peaks, SO, SA and the volume mode (V) are identified based on Maxwell and Landau-Lifshitz equations mentioned in previous section. The spectrum in Fig. 4.8a corresponds to the dynamic magnetization amplitudes depicted in Fig. 4.3c and d with parallel alignment of $M_1$ and $M_2$ (magnetizations of the two permalloy layers). Here the SW amplitude of the SA (SO) mode is suppressed (enhanced) from its zero $I_{channel}$ value (Fig. 4.3a and b) in the
left ferromagnetic layer on which the probe laser beam is incident. In the spectrum of Fig. 4.8b, $H_{\text{cur}} (I_{\text{channel}})$ is reversed. This corresponds to Fig. 4.3e and f, where now the SA (SO) amplitude is instead enhanced (suppressed) from its zero $I_{\text{channel}}$ value. These changes to the SW amplitudes in Fig. 4.3c to f are consistent with the observed changes of SA to SO BLS peak ratio from Fig. 4.8a (91/78=1.1) to Fig. 4.8b (128/27=4.7). Fig. 4.8c and 4.7d present similar spectra at a much higher external field ($H_{\text{ext}}=-5.19\text{kOe}$) and the SO/SA peak intensity switching feature complies with those of Fig. 4.8a and 4.8b. Fig. 4.8e illustrates a BLS spectrum at a regime where $|H_{\text{cur}}| > |H_{\text{ext}}|$. In this case the Stokes surface SW frequencies (Fig. 4.5b) do not match their anti-Stokes counterpart confirming anti-parallel alignment of $M_1$ and $M_2$ as in Fig. 4.6a and b. Finally when $H_{\text{ext}}$ is reversed with $H_{\text{cur}}$ unchanged, the intensity switching evident in Fig. 4.8f corresponds to the case in Fig. 4.6e and f.
Figure 4.8: Anti-Stokes Brillouin spectra recorded in sample #2 for various values of $H_{\text{ext}}$ and $H_{\text{cur}}$ ($I_{\text{channel}}$). The numbers beside each peak indicate the peak intensity (subtracted from background). All spectra are recorded with light incident on same magnetic layer.

In order to quantify the dependence of the mode intensity switching on the effective magnetic field, we define the switching strength parameter $S$ at each value of $|H_{\text{ext}}|$ and $|H_{\text{cur}}|$, given by $S(|H_{\text{ext}}|,|H_{\text{cur}}|) = R(\text{same}) / R(\text{opposite})$ [44]. Here $R(\text{same})$ and $R(\text{opposite})$ are the peak intensity ratios of the higher ($\omega_h$) to lower ($\omega_l$) frequency.
surface SW in the BLS spectrum when $H_{\text{ext}}$ and $H_{\text{cur}}$ are respectively directed in the same and opposite directions on the magnetic layer on which the probe laser beam is incident. Fig. 4.9 summarizes the measured data (from both sample #1 and #2) with calculated fits based on Maxwell and Landau-Lifshitz equations of the switching strength $S$ as a function of $|H_{\text{ext}}|$ for different values of $|H_{\text{cur}}|$. The calculation adopted $M_0$ (saturation magnetization) = 774 Oe, $d = 32.2$ nm and $d_0 = 45$ nm to achieve the best agreement with observed SW frequencies and the parameter $S$. It is noted from Fig. 4.9 that for antiparallel alignment ($|H_{\text{ext}}| < |H_{\text{cur}}|$) of $M_1$ and $M_2$, calculations show that as $|H_{\text{ext}}|$ increases $S$ rises on a single curve for different Amperian fields $|H_{\text{cur}}|$. On the other hand, as $|H_{\text{ext}}|$ is swept beyond this range, $S$ sequentially collapses to values corresponding to parallel alignment when $|H_{\text{ext}}|$ surpasses $|H_{\text{cur}}|$. The switching ratio now lies on separate curves for different $|H_{\text{cur}}|$ with the curves remaining relatively flat for $|H_{\text{ext}}| < 1$ kOe and subsequently rising gradually with $|H_{\text{ext}}|$ beyond 1 kOe (curves with larger $|H_{\text{cur}}|$ display greater rate of increase). These trends of the dependence of $S$ on $|H_{\text{ext}}|$ and $|H_{\text{cur}}|$ are also reflected fairly well in the experimental data (Fig. 4.9).
Figure 4.9: Solid lines represent calculated switching strength parameter $S$ as a function of $|H_{\text{ext}}|$ for various values of $|H_{\text{cur}}|$. An experimental point is hollow (filled) if spectra correspond to anti-parallel (parallel) alignment of magnetization is detected.

4.5 Intuitive analysis

In order to gain physical insight into the results summarized in Fig. 4.9 we introduce a mechanical analog of $M_1$ and $M_2$, namely a coupled pendulum (spring constant $k$) (Fig. 4.10a) where the two masses ($m$) experience different gravitational fields $g+\Delta g$ and $g-\Delta g$. The equations of motion are given by

$$
\begin{align*}
mx_1 + (k + \frac{m(g + \Delta g)}{l})x_1 - kx_2 &= 0 \\
mx_2 + (k + \frac{m(g - \Delta g)}{l})x_2 - kx_1 &= 0
\end{align*}
$$
Setting $\alpha = \frac{k}{m}$ and $l = 1$ and seeking solutions in the form $x_\sigma = \rho_\sigma e^{i\omega t}$, the eigenvalues and eigenvectors for the pendulum normal mode are respectively

$$\omega^2 = g + \alpha \pm \sqrt{\alpha^2 + \Delta g^2}$$

and

$$\begin{bmatrix} \rho_1 \\ \rho_2 \end{bmatrix} = \frac{1}{\sqrt{1 + f_\pm^2(\eta)}} \begin{bmatrix} f_\pm(\eta) \\ 1 \end{bmatrix}$$

where $f_\pm = -\eta \pm \sqrt{1 + \eta^2}$ and $\eta = \Delta g/\alpha$. Fig. 4.10b and c show the dependence of $\rho_1$ and $\rho_2$ on $\eta$ for the low (l) and high (h) frequency normal modes respectively. It is noted that the eigenvectors solely depend on the ratio $\Delta g/\alpha (= \eta)$, and are independent of $g$.

We are now ready to explain the change in SO/SA spin wave precession amplitudes depicted in Fig. 4.3 using this coupled pendulum model. In Fig. 4.10b and c, when $\eta = 0$, the amplitudes of the two pendulums are identical for both the lower frequency (in-phase) mode and higher frequency (out-of-phase) mode. This corresponds to the equality of spinwave amplitude on the two magnetic layers for both the SA spinwave and SO spinwave in Fig. 4.3a and b. In Fig. 4.3c and d, a field imbalance is turned on such that the effective field on the left magnetic layer is reduced to $H_{\text{ext}} - H_{\text{cur}}$ while the effective field on the right layer is enhanced to $H_{\text{ext}} + H_{\text{cur}}$. This corresponds to a negative value of $\eta$ in Fig. 4.10a. In Fig. 4.10b, for a negative value of $\eta$, the amplitude of the lower frequency normal mode is enhanced (reduced) on the left (right) hand side pendulum. This explains the amplitude modification of the SO spinwave depicted in Fig. 4.3d recalling that the SO spinwave is the lower frequency mode (Fig. 4.2). Notice that the coupled pendulum is energetically lower for in-phase oscillation than out-of-phase, which is opposite to spin wave oscillation in coupled parallel-aligned magnetic bilayers.
Magnetization on bilayers is energetically costly in parallel-alignment and therefore produces a lower frequency of oscillation for out-of-phase (SO) spin wave mode than in-phase (SA) mode. Similarly, for a negative value of $\eta$, the amplitude of the higher frequency normal mode is reduced (enhanced) on the left (right) hand side pendulum (Fig. 4.10c). This corresponds to the amplitude modification of SA spinwave depicted in Fig. 4.3c given that SA spinwave is the higher frequency mode (Fig. 4.2). Precession profile of SA and SO mode depicted in Fig. 4.3e and f can be explained in the same manners with $\eta$ reverses to positive value.

Returning to addressing the results summarized in Fig. 4.9: In an anti-parallel alignment configuration, $H_{\text{cur}}$ and $H_{\text{ext}}$ are equivalent to $g$ and $\Delta g$ respectively. Therefore, in analogy to the pendulum eigenvectors, modifications to the SW mode amplitudes, and thus the parameter S, are independent of $H_{\text{cur}}$ and depend only on $H_{\text{ext}}$. This accounts for the rise of S on a single curve for different $H_{\text{cur}}$ in Fig. 4.9 in the anti-parallel region. On the other hand, since $H_{\text{ext}}$ and $H_{\text{cur}}$ switch roles in being represented by $g$ and $\Delta g$ when transiting from anti-parallel to parallel alignment, after the transition S does not depend on $H_{\text{ext}}$ (equivalent to $g$) anymore but is now sensitive only to $H_{\text{cur}}$ ($\Delta g$). Therefore, in Fig. 4.9, after entering parallel alignment (for $H_{\text{ext}} < 1\text{kOe}$) S transfers to a distinct set of flat curves for different $H_{\text{cur}}$. The gradual increase of S beyond $H_{\text{ext}} > 1\text{kOe}$ in Fig. 4.9 is explained in terms of the frequency splitting between the SA and SO modes. Fig. 4.8 shows that the frequency splitting decreases from 2.9GHz (Fig. 4.8a and b) to 1.5GHz (Fig. 4.8c and d) as $|H_{\text{ext}}|$ increases from 0.54kOe to 5.19kOe. A decrease in $(\omega_h - \omega_l)$ indicates a drop in the value of $\alpha$ which results in larger $\eta (=\Delta g/\alpha)$ for a given value of $\Delta g$. Fig. 4.10b and c illustrates that a larger $\eta$ corresponds to larger deviation of $\rho_{hi}(2)$ and
\( \rho_{l(2)} \) from their \( \eta = 0 \) values which would lead to a larger switching strength \( S \). On the other hand, there exists a discontinuous jump of the surface SW frequency splitting that accompanies the anti-parallel to parallel transition (0.9GHz of Fig. 4.8e and f to 2.9GHz of Fig. 4.8a and b). This explains the abrupt drop of \( S \) across the anti-parallel/parallel boundary in Fig. 4.9.

Figure 4.10: (a) Coupled pendulums with different local gravitational fields \((g+\Delta g)\) and \((g-\Delta g)\). (b) and (c) illustrate dependence of normalized eigenvector elements \( \rho_1 \) (solid line) and \( \rho_2 \) (dash line) on \( \eta (=\Delta g/\alpha) \) for the lower (l) and higher (h) frequency normal modes.
4.6 Conclusion

Interesting results associated with “cross-communication” between the normal modes emerge when current pulses, rather than continuous channel currents are utilized. In the pendulum analogy, it can be shown that when the constant $\Delta g$ in Fig. 4.10 is replaced by a field pulse of strength $\Delta g$ and time duration $\Delta t$, the amplitude of the higher and lower frequency eigen-modes after passage of the pulse are, in general, a linear combination of the amplitudes of these two modes prior to the pulse. However for $\Delta g >> \alpha (\eta >> 1)$ if the duration and height of the pulse is tailored to satisfy: 1) $\Delta t = \pi/(\omega_h - \omega_l)$ and 2) $\omega_l = [n/(n+1)]\omega_h$ [45], where $\omega_h (\omega_l)$ are the frequency of higher (lower) modes during application of the $\Delta g$ pulse and $n$ is an integer, then the lower and higher frequency mode will neatly exchange their amplitudes (instead of remaining in the aforementioned mixed state) after passage of the $\Delta g$ pulse, thus achieving cross-communication between vibrational eigen-modes. Given that the eigenvectors do not depend on $g$, it follows that similar cross-communication between modes will not occur if a $g$-pulse replaces the $\Delta g$-pulse. Therefore, this mode cross-communication represents a dynamic mechanism that is unique to the coupled SW system with the freedom to tune the field imbalance.

In summary, we have investigated through Brillouin light scattering the effects of magnetic field imbalances on the precessional amplitudes of SWs in a magnetic double layer. The imbalance, generated by an Amperian field and an external field, allows for the eigen-modes of the precessing magnetizations to be locally tuned. The results are explained by conventional electrodynamics and within a mechanical coupled pendulum analog in an unequal gravitational field. Transfer of amplitude information from one
mode to another by pulsed fields (equivalent to $\Delta g$-pulses), i.e. the cross-communication between SWs, is shown to be a distinctive feature of the tunable unbalanced local fields. It should be noticed that the physical system studied (both experimentally and theoretically) in this chapter represents a very general physical circumstance. Although it is a two magnetic slabs system which is investigated, the same physics should apply even to two single classical spins. The coupling between the two magnetic layer are purely magnetostatic, but replacing it with exchange coupling will not change the underlying physics. Also origin of field imbalance are not limited to combination of $H_{\text{ext}}$ and $H_{\text{cur}}$ as depicted in this chapter, but can arise due to difference in local crystal strain (magnetostriction) experienced by coupled spins.
CHAPTER 5

SPIN-POLARIZED CURRENT AND DOMAIN WALL MOTION

5.1 Introduction

The discovery of giant magnetoresistance (GMR) [1] in 1988 illustrated that the relative orientation of magnetization in multilayers controls the resistance and, thus, influences the local electric current flowing through multilayer structures. However the reverse effect: namely, the effect of current on the orientation of the magnetization was only predicted [46] and experimentally observed [47] in the 1990’s. It was envisioned and experimentally proven [46,47] that, if the electrical current density is high enough, the spin polarized characteristic of the current will stimulate oscillation [46,47] in the local magnetization or even reverse the magnetization [46,47]. This phenomena, named “spin torque transfer”, is schematically illustrated in Figure 5.1.

In Fig. 5.1a, a metallic magnetic micro-wire is displayed. The wire contains two magnetic domains with exactly opposite orientation of magnetization. The boundary, or the domain wall (DW), is located at the center. When a voltage is applied to the wire, a current of electrons will flow through the wire (current with a particular spin is called a “spin-polarized” current). In the figure, a spin-up electron (same direction as the local magnetization) moves from right to left. In a magnetic material, scattering of electrons
is spin-dependent. Electrons with spin oriented opposite to the local magnetization undergo more scattering than electrons with spin parallel to the local magnetization. Thus electrons whose spin are aligned with the local magnetization has a lower effective electrical resistance.

In Fig. 5.1b, the spin up electron has moved across the boundary (domain wall) and entered a region with magnetization pointing downward. As illustrated in Fig. 5.1c, in order for the electron to experience a lower resistance and continue moving to the left it has to flip its spin lie along the magnetization of the left domain. In order to conserve angular momentum, the local magnetization acquires the angular momentum change of the electron. Fig. 5.1d illustrates this reversal of the local magnetic moment thereby moving the domain wall one local moment to the left. The electron, with its own spin flipped, will continue to move to the left until it reaches the end of the wire (Fig. 5.1d).

As additional electrons continue to move from right to left, the domain wall will be pushed further towards the left. The net result is the movement of the domain wall which is directly controlled by the spin-polarized electric current. This is a simplified 1-D view of the spin torque transfer mechanism. In reality the domain wall has a finite width that can lead to deviation from this simplified model.

As originally proposed by Berger, such a spin torque transfer mechanism can also be understood from the point of view of “spin accumulation” [48]. Spin-up electrons from the right are continuously injected across the domain wall to produce an excess spin-up electron population over spin-down. This population imbalance leads to an
energy gap between the Fermi surfaces of the two spin bands. The spin torque transfer can be viewed as magnon (spin-wave) emission due to electrons hopping from one spin band to the other.

Figure 5.1: Illustration of spin-torque transfer mechanism.
The spin torque transfer was experimentally first investigated in the current-perpendicular-to-plane (CPP) geometry [47]. In recent years, it was realized that the same mechanism is also applicable in the current-in-plane (CIP) direction. Magnetic domain walls (DW) in micro- and nano-wires with current flow serves as an excellent system for the CIP spin-torque transfer phenomenon to occur. Clear observations of spin torque in permalloy nanonwires were carried out by magnetic force microscopy (MFM) [49, 50], and spin electron microscope (SEM) [51]. In these experiments, both head-to-head and tail-to-tail DWs move opposite to the direction of the applied current (i.e. same direction as electrons flow). This is a characteristic spin-torque feature distinguishing the phenomenon from the effect associated with the Amperian magnetic field of the current. If the DW motion arises from the Amperian field, the two types of domain walls should move along opposite directions in response to the same applied electric current.

When magnetic elements get smaller in dimensions, the spin torque transfer mechanism is preferred over exploiting Amperian field to induce magnetization reversal. There are a couple of factors that make the Amperian field ineffective in reversing nanoscale magnetic elements. First the maximum current and the associated Amperian field that can be carried by such miniature elements decreases with its size. Moreover in order to overcome super-paramagnetism [52], the magnetic element must be increasingly coercive. On the other hand, since the spin torque transfer mechanism is proportional to the current density (and not the magnitude of the current), shrinkage of the magnetic
element poses no burden on the effectiveness of the spin torque transfer mechanism. Further, the spin torque transfer mechanism in a CIP geometry can also reduce crosstalk between adjacent magnetic wires.

In this chapter the CIP spin torque transfer phenomena in fabricated micro-wires observed by Kerr imaging is presented.

5.2 Methodology

Figure 5.2 displays a sample utilized in our spin-torque experiment. A permalloy film of thickness 50 nm was deposited on SiO$_2$(100 nm)/Si substrate by means of magnetron sputtering. The film was then patterned by photolithography and wet-etched into the channel structure shown in Figure 5.2. The central channel is 125 µm long and 8 µm wide. The long axis of the channel was chosen to lie parallel to the intrinsic anisotropy easy direction of the film. As shown the channel is connected to the film on one side by a tapered structure and to the other by a series of micro-wires.

During the experiment, the sample is electrically connected to an external circuit as shown in Figure 5.3. Approximately each second the relay switch reverses connection. When linked to the left, the dc supply between 0 and 100 V charges the capacitor through the large resistor R that restricts the current drawn to be small. At the same time R is small enough that RC (time constant) is much smaller than one second. When the relay switch links to the right, the capacitor discharges through the sample. Given that the capacitance is 2.5 nF, and the total sample resistance is 220 Ω, the discharge time constant is approximately half a micro-second.
Figure 5.2: Image of channel utilized. (a) A wide field view of the structure. The central channel is connected to large area of film by tapered structure on one side, and by a series of micro-wires on the other side. (b) A higher magnification view of the central channel. The dimension in microns of the channel width/length = 8/125.
5.3 Experimental results

The magnetization of the channel was monitored using Kerr microscopy as described in chapter 2. A magnetic field of several hundred Gauss was applied along the long axis of the channel to magnetically saturate the sample. Subsequently, the external magnetic field was reduced to zero and a reversed field of several Gauss applied. In such a situation, the permalloy at the unpatterned film area on the two ends of the sample (Fig. 5.2a) will have their magnetization completely reversed. However, since within the channel and the outer micro-wires, a reversal field of several Gauss is not strong enough to produce domain nucleation, their respective magnetization remains unchanged. If the reverse field is increased incrementally, domains existing in the two unpatterned film regions at the sample edges will enter the central channel via domain wall motion. However, the wall movement from each end does not travel symmetrically towards the channel. The wall moving from the left along the tapered structure (Figure 5.2a), encounters less resistance than the wall moving through the micro-wire structures
that have a greater number of pinning centers. Therefore a domain wall from the left always reaches the channel first thus ensuring the existence of initially only one (and not two) DW’s entering the central channel.

After the domain wall from the left enters the channel, the external field is reduced to zero. The corresponding Kerr image is shown in Figure 5.4a. Figures 5.4(b) to 5.4(i) show the subsequent evolution of the domain wall. During the time between taking each image (Fig. 5.4), two current pulses [53] are applied through the channel according to the methodology described in section 5.2. The arrow next to each image indicates the direction of electron flow of the applied current pulse before the image is taken. As seen from Fig. 5.4b, after application of the first two pulses, the DW moves many micrometers towards the center. Fig. 5.4c shows the retraction of the DW to the left after 2 current pulses in the reverse direction. The back and forth movements of the DW with current in Fig. 5.4d and e is obvious. Subsequent images (Fig. 5.4e to i) however shows a weaker DW response to current pulses.
Figure 5.4: Response (Kerr images) of a head-to-head domain wall to electrical current pulses. Arrows indicate direction of electron flow in the current pulse before each image is taken.

Figure 5.5 shows the response of a tail-to-tail domain wall to current pulses. The initial position of the domain wall within the channel is displayed in Fig. 5.5a. The back and forth movements of the domain wall after application of current pulses in opposing directions in Fig. 5.5b to 5.5e are obvious. In later figures, i.e. Fig. 5.5f and g, the responses become less obvious.
Figure 5.5: Response (Kerr image) of a tail-to-tail domain wall to electrical current pulses. Arrows indicate direction of electron flow in current pulse before each image is taken.

Figure 5.6 shows a series of Kerr images also taken on a tail-to-tail domain wall, in another sample with structural dimensions identical to the sample illustrated in Fig. 5.4 and 5.5. We observe in Fig. 5.6(a) to (d) the wall movement is correlated with the direction of electron movement. Clear movement of the domain wall is evident after every two pulses in each image. The movement in Fig. 5.6d is large and is accompanied by the break up of the domain wall into segments.
Figure 5.6: (a) to (d): Example of the response of a tail-to-tail domain wall to electric current from a sample with identical pattern dimension as previous samples (Fig. 5.4, 5.5). Arrows indicate the direction of electron flow in current pulse before each image is taken.

(e): Depiction of segmental domains inside the channel.

Domain wall movement is only observed when the applied dc voltage is higher than 60 Volt. The images shown in Figure 5.3, 5.4 and 5.5 have the dc voltage set at 72 Volt. Taking account of the width and thickness of the channel, this corresponds to a current density of $8.2 \times 10^{11} \text{Am}^{-2}$.

Overall, the data displayed in figures 5.4 to 5.6 confirms that the domain wall always moves in the same direction as electron flow, regardless of whether the DW is head-to-head or tail-to-tail. This is a significant characteristic, indicating that the movement is produced by the spin-torque transfer mechanism and not by the
Amperian field associated with the current pulse. If the response is due to the Amperian field, the direction of movement of head-to-head DW would be opposite to that of the tail-to-tail DW for similar current pulses.

Although the data in Fig. 5.4 to 5.6 do reveal the direction of DW movement to be correlated with the direction of electron flow, there were occasional situations where the DW movement was opposite to the electron flow direction. We have eliminated the possibility of interference of Amperian fields from the external wires needed for the circuitry. The channel was also aligned in the east-west direction such that Earth’s magnetic field would not influence the response. We believe the occasional inconsistency of DW movement with the pure spin-torque effect is related to the width of the channel. The channel is 8 µm wide and could allow the domain wall to form a complex structure, thereby making the exact location of the domain wall not well defined. The domain wall can always rearrange itself after responding to the current pulses due to the dipolar-dipolar micro-magnetic interactions. This wall rearrangement can obscure the initial response of the domain wall to the current and lead to a response opposite to what is expected solely from the spin-torque mechanism. The same mechanism may also lead to breakdown of the domain into segments as illustrated in Fig. 5.6 a to d, with domain patterns inside the channel as depicted in Fig. 5.6(e). The diminished wall movement as more current pulses are applied (Fig. 5.4e - i and Fig. 5.5f, g) may arise from the fact that magnetic domains within the channel have converged into certain metastable configurations due to magnetostatic interactions. On the other hand, a strong repulsive force on the domain wall from the tapered structure on the extremity of the channel has been observed (domain wall near to the channel extremity always moves toward the...
central part of the channel regardless of applied current direction), which agrees with micromagnetic simulation. This supports that, in our relatively wide channels, magnetic dipolar-dipolar forces do influence behavior of the domain wall. Narrow wires with width equal to 1 µm were fabricated but collecting Kerr images on these wires were not successful during the first several attempts.

5.4 Computer simulations

5.4.1 One dimensional

In this section we investigate, through simulations, the physics underlying the spin-torque transfer mechanism that leads to domain wall motion. It has been known that in the current-perpendicular to plane (CPP) geometry, a dc current generates high frequency magnetization oscillation by the spin-torque transfer mechanism [10]. It has been speculated that a similar effect can occur in the CIP geometry, when current traverses a DW [49].

Let us first consider a one dimensional model. A series of equally spaced spins is placed along the x axis, as illustrated in Figure 5.7a. The dynamics of \( i \)th spin is governed by the Landau-Lifshitz-Gilbert equation (with spin torque included):

\[
\dot{S}_i = \gamma_0 H_i \times S_i + \alpha \gamma_0 S_i \times (H_i \times S_i) + \kappa (S_{i+1} - S_i)
\]

(5.1)

where \( S_i \) is a unit vector. In equation (5.1), \( \gamma_0 = \mu_0 g / |\mu_B| / h \) and \( \kappa = JgP \mu_B / 2eMs \Delta x \), where \( \Delta x \) is the spacing between spins, \( J \) is the electrical current density and \( P \) is the spin-polarization rate \( (P = (j^\uparrow - j^\downarrow) / (j^\uparrow + j^\downarrow)) \), where \( j^\uparrow \) and \( j^\downarrow \) are the current density due to spin-up and spin-down electrons). The 1st term on the right side of (5.1) represents
precession, while the 2\textsuperscript{nd} term gives rise to damping since this term always tends to align $\mathbf{S}_i$ with the effective field $\mathbf{H}_i$. The 3\textsuperscript{rd} term represents the spin torque transfer effect. Since the exact form of this term is still in dispute [54], we follow the model of Ref. 54 and adopt the simplest form. While $\mu_0$ (permeability constant), $g$ (Landé g-factor), $\mu_B$ (Bohr magneton), $h$ (Planck constant), $e$ (electron charge) take their standard numerical values, we adopt $\alpha$ (Gilbert damping factor)=0.02, $P$ (electron polarization rate)=1 and $M_s$ (saturation magnetization)=8×10\textsuperscript{5} Am\textsuperscript{-1}.

The effective magnetic field experienced by the $i$\textsuperscript{th} spin, written as $\mathbf{H}_i$, is the summation of exchange field ($\mathbf{H}_{\text{exc,}i}$), anisotropy field ($\mathbf{H}_{\text{ani,}i}$), magnetic dipolar interaction field ($\mathbf{H}_{\text{dip,}i}$) and external field ($\mathbf{H}_{\text{ext,}i}$):

$$\mathbf{H}_i = \mathbf{H}_{\text{exc,}i} + \mathbf{H}_{\text{ani,}i} + \mathbf{H}_{\text{dip,}i} + \mathbf{H}_{\text{ext,}i}$$

(5.2)

The exchange field is usually related to the 2\textsuperscript{nd} derivative of the magnetization by [55]

$$\mathbf{H}_{\text{exc}} = \frac{2A}{M_s} \nabla^2 \mathbf{M}$$

(5.3)

where $A$ is the exchange stiffness constant. The value of $A$ for Permalloy equals 13 pJ/m.

In discrete form,

$$\nabla^2 \mathbf{M} = \frac{\mathbf{M}_{i+1} + \mathbf{M}_{i-1} - 2\mathbf{M}_i}{\Delta x^2} = \frac{M_s (\mathbf{S}_{i+1} + \mathbf{S}_{i-1} - 2\mathbf{S}_i)}{\Delta x^2}$$

(5.4)

Taking account of the fact that the cross product of $\mathbf{S}_i$ with itself vanishes, the exchange field $\mathbf{H}_{\text{exc,}i}$ experienced by the $i$\textsuperscript{th} spin can therefore be expressed as:

$$\mathbf{H}_{\text{exc,}i} = \frac{2A}{M_s \Delta x^2} (\mathbf{S}_{i+1} + \mathbf{S}_{i-1})$$

(5.5)
The anisotropy field experienced by the ith spin is given by the dot product of the fixed anisotropy field $H_{ani}$ with $S_i$:

$$H_{ani,i} = \left( \frac{H_{ani} \cdot S_i}{|H_{ani}|} \right) H_{ani}$$

(5.6)

In the one dimensional case, we replace the dipolar interaction field $H_{dip,i}$ by a simple demagnetization field [55]:

$$H_{dem,i} = -M_x (S_i \cdot z)$$

(5.7)

A LabView program (available in electronic form upon request [56]) was written to implement this one-dimensional model. Spins are placed along the x axis and their orientation initially aligned along the x axis, as shown in Figure 5.7a. A domain boundary exists at the center. In this one dimensional simulation, the exchange interaction between adjacent spins is initially turned off by assigning the exchange stiffness constant $A$ to zero. The effect of $A \neq 0$ is discussed later. The electrons flow along the negative x direction and the current density is set at $14 \times 10^{11}$Am$^{-2}$. In the simulation, the first spin encountered by the electron after passing the domain boundary was monitored in real time. The time dependence (in picoseconds) of the x, y and z components of this spin is plotted in Fig. 5.7b, c and d. The average value of the spin components against time is always subtracted from the spin components themselves to enable the LabVIEW diagram to be sensitive to small oscillations. Therefore only the relative values of the spin components (and not their absolute values) are considered. As seen from these figures, the current generates a high frequency oscillation (roughly at 30 GHz) of the three spin components. The x, y and z oscillation amplitudes amplify until they reach a saturation value beyond which a stable oscillation is sustained. Although rapid oscillations occur,
the spin direction does not flip as the x component of the oscillation spin amplitude is 0.55, which is less than 1 (Fig. 5.7b) (since \( S_x = -1 \) initially, it requires a relative change of spin x component to equals 1 in order to bring \( S_x \) to a positive value). Fig. 5.7e shows the x component of the monitored spin if a higher current density of \( 20 \times 10^{11} \text{Am}^{-2} \) is used instead. In this case the x component of the spin transited from -0.75 to 1.25, corresponding to a total transition magnitude of 2, indicating that the spin has flipped from -1 to +1 and has thus completely switched direction.

The situation becomes very different if the exchange interaction \( A \) is turned on during the oscillation. The evolution of the spin components against time is shown in Fig. 5.8. In the beginning, \( A \) is set to zero and a current density of \( 14 \times 10^{11} \text{Am}^{-2} \) with electron flow in the negative x direction. The y and z spin components oscillate with increasing amplitude. If this amplification of the oscillation amplitudes continue, they will reach the same sustained stable oscillation as shown previously in Figure 5.7. However, at \( t=170 \text{ ps} \), the exchange stiffness constant \( A \) is abruptly introduced and increases from zero to 0.3 pJ/m. As evident from Fig. 5.8 b and c, the oscillation is subsequently attenuated. It appears that the exchange interaction has a detrimental effect on magnetization oscillation. In fact if \( A \) increases further, spin simply switch direction from \((S_x, S_y, S_z) = (-1,0,0)\) to \((S_x, S_y, S_z) = (1,0,0)\) smoothly without any oscillation in between. Oscillatory solution were not found for \( A \) larger than 0.3 pJ/m.

In conclusion, in the one dimensional spin torque model, while dc currents produce high frequency oscillation in the CPP geometry, the same mechanism is not effective within a domain wall in the CIP geometry for finite exchange interaction between adjacent magnetic moments in the domain wall with finite width. In the CPP
geometry, adjacent magnetic moments along the electron path reside on different magnetic layers and are separated from each other by a relatively thick non-magnetic metallic spacer. Therefore the spins are decoupled from each other. In contrast, in the CIP geometry, adjacent magnetic moments are strongly exchange coupled. Such coupling damps any oscillation directly generated by the spin-torque transfer mechanism. Therefore, any spin wave stimulation in 1D that may occur in the CIP spin-torque transfer experiment must be a higher order effect and does not have the same origin as their CPP counterpart. Although this simulation was carried out for a one-dimensional (1D) case, similar conclusions should also apply for real physical magnetic wire which are in essence two dimensional (2D) The additional degree of freedom in a 2D system do not provide a mechanism to turn off the strong coupling between neighboring local magnetic moment.
Figure 5.7: (a) Spins are aligned along the x axis. A domain wall is placed at the center. Exchange interaction is turned off (A=0). Since there is no exchange interaction, the domain wall width is zero and we have an abrupt wall. (b) x component of the monitored spin at domain wall. (c) y component of the monitored spin. (d) z component of the monitored spin. Because of demagnetization field (equation (5.7)), oscillation amplitude in z direction is much smaller than in y. Visualization of the spin oscillation is given in the inset of (a). Only the relative (not absolute) values of the amplitude components matter and are shown in (b)—(e). Current density is set at $14 \times 10^{11} \text{Am}^{-2}$. (e) x component of the monitored spin against time if a current density of $20 \times 10^{11} \text{Am}^{-2}$ is used instead. The unit for time is picosecond. Vertical axes in (b) to (e) correspond to relative spin component oscillation amplitude and are unit-less.
Figure 5.8: Amplitude of the monitored spin (a) x component, (b) y component and (c) z component. Oscillation amplitude in x direction remains nearly zero since by comparison (to Fig. 5.7) the amplitudes in y and z direction are extremely small. Initially, a current of $14 \times 10^{11} \text{Am}^{-2}$ is applied with the electron flow along -x and exchange stiffness $A=0$. At $t=170 \text{ ps}$ in the figure, $A$ abruptly increases to 0.3pJ/m. Time $t$ is in picoseconds.

5.4.2 Two dimensional

One of the major deficiencies in the one dimensional model discussed in the previous section is the inadequate treatment of magnetic dipole-dipole interactions. This interaction must be addressed in a two dimensional regime.
The magnetic field at \( \mathbf{r} \) due to a magnetic moment \( \mathbf{m} \) at the origin can be expressed as [16]:

\[
\mathbf{H} = \frac{3(\mathbf{m} \cdot \hat{\mathbf{r}}) \hat{\mathbf{r}} - \mathbf{m}}{4\pi r^3}
\]  

(5.8)

where \( \hat{\mathbf{r}} \) indicates the unit vector of \( \mathbf{r} \). Writing the magnetic moment at location \( \mathbf{r} = (i\hat{x} + j\hat{y})L \) as \( \mathbf{m}(i,j) = M_s(i,j) \), where \( \mathbf{S} = S_x \hat{x} + S_y \hat{y} + S_z \hat{z} \), \( i, j \) are integers, \( L \) is the grid size and \( L_z \) is the magnetic film thickness. It can be derived straightforwardly that the magnetic field at location \( \mathbf{r} = (p\hat{x} + q\hat{y})L \), where \( p, q \) are integers, due to \( \mathbf{m}(i,j) \) is given by (see Figure 5.9)

\[
\mathbf{H}(p-i, q-j) = \frac{L_z M_s}{4\pi L} \left[ \left\{ N_{xx}(p-i, q-j)S_x(i, j) + N_{xy}(p-i, q-j)S_y(i, j) \right\} \hat{x} + \left\{ N_{xy}(p-i, q-j)S_x(i, j) + N_{yy}(p-i, q-j)S_y(i, j) \right\} \hat{y} + N_{zz}(p-i, q-j)S_z(i, j) \hat{z} \right]
\]  

(5.9)

Where the dipolar interaction tensors,

\[
N_{xx}(p,q) = \frac{2p^2 - q^2}{(p^2 + q^2)^{3/2}}
\]  

(5.10a)

\[
N_{xy}(p,q) = \frac{3pq}{(p^2 + q^2)^{3/2}}
\]  

(5.10b)

\[
N_{yy}(p,q) = \frac{2q^2 - p^2}{(p^2 + q^2)^{3/2}}
\]  

(5.10c)

\[
N_{zz}(p,q) = \frac{-1}{(p^2 + q^2)^{3/2}}
\]  

(5.10d)

have purely geometrical origin and do not depend on the local spin \( \mathbf{S}(i,j) \). The total magnetic field at \( \mathbf{r} = (p\hat{x} + q\hat{y})L \) due to all magnetic moments at different locations is given by
\[ H_{p,q} = \sum_{i,j} H(p-i,q-j) \]  

(5.11)

As seen from (5.9) and (5.11), calculation of the dipolar magnetic field involves a two-dimensional convolution of spin components \( S_x, S_y \) and \( S_z \) with the dipolar interaction tensors \( N_{xx}, N_{xy}, N_{yy} \) and \( N_{zz} \). We take advantage of two-dimensional fast Fourier transform [57] to accelerate the computation. The convolution in real space becomes simple multiplication in reciprocal space, making computation in reciprocal space more favorable.

Figure 5.9: Coordinate system of the magnetic moment \( \mathbf{m} \). \( L \) is the grid size utilized in the simulation.

A Visual C++ program (2DLLES [56]) was written to simulate the behavior of the domain wall inside the channel in the two-dimensional case. By relaxing the magnetization inside the channel with appropriate predefined orientation, two types of domain walls are found to be stable in a permalloy nanowire, which are the “Vortex” wall (Fig. 5.10a) and “Transverse” wall (Fig. 5.10b). The width of the wire is set to be 250 nm and thickness equal 25 nm. The “Object Oriented MicroMagnetic Framework”
(OOMMF) [58], developed at the National Institute of Standard and Technology (NIST) also gives the same result. The response of these two types of domain walls with respect to external magnetic field and spin polarized current were probed. It was found that the results depend largely on chosen grid size during the simulation. While similar results were found in narrower wires, wires with larger width are not simulated because of their required longer computational time.

![Figure 5.10: Result of micromagnetic simulation. (a) Vortex domain wall (b) Transverse domain wall](image)

5.5 Conclusion

The subject of domain wall motion induced directly by a spin-polarized current is presently an active research area in condensed matter physics. The experiments in Ref. 49 to 51 show clearly the effect in submicron wide permalloy wires. In these experiments, magnetic force microscopy and spin-polarized scanning electron microscopy were used to monitor the domain walls. In the wires with submicron width utilized in these
experiments, domain walls are better defined and observations of the effect are conclusive. On the other hand, it is demonstrated in our experiments that wide field Kerr imaging techniques is capable of revealing the spin-torque transfer effect in a CIP geometry and the resulting DW motion was observed. The 8 µm wide wires utilized in our studies however enable the domain walls to rearrange themselves when the current is reduced to zero. Domain walls residing in the wider wires of our studies suffer disturbances from magnetostatic interaction within the wires. In the future, thinner wires (with thickness equal 10 nm, instead of our present value of 50 nm) should be used instead, while the wire width can be kept the same (8 µm). This will suppress undesired magnetic dipole-dipole interactions. In such circumstances, substrates with better surface quality (smoother) should be used since thinner wires are more susceptible to pinning effects arising from surface roughness. The substrate surface involved in our studies was generally too rough and 10 nm thick fabricated wires suffered strong pinning effects. Experiments on these thinner wires should only be done when the substrate surface condition is improved.

Another possible future direction lies in the usage of synchrotron radiation facilities at Argonne national laboratory [59]. With the much shorter wavelength of x-rays, better spatial resolution will be provide and probing of sub-micron magnetic wires will become possible. In these narrower wires, domain walls are better defined and are free from all magnetostatic interactions that adversely affect the spin torque response.
CHAPTER 6

MAGNETIZATION REVERSAL IN CO/PT MULTILAYERS

6.1 Introduction

With rapid advances taking place in the magnetic recording industry, understanding magnetization reversal processes, including the formation and evolution of magnetic domains, have emerged as important issues in data storage and sensor technology. As current magnetic storage density is pushed towards the superparamagnetic limit, it is believed that higher areal densities can be achieved by means of perpendicular recording due to the lower superparamagnetic limit and larger anisotropy in magnetic media with perpendicular anisotropy [60]. Because there are many competing energies involved in the magnetic switching processes, the formation and evolution of magnetic domains during magnetization reversal can be complex and fundamentally interesting.

The impressive progress in data storage science and development over the past two decades has been driven, in part, by discoveries of oscillatory interlayer coupling in ferromagnetic (FM) and nonmagnetic (NM) multilayers and the occurrence of giant magneto resistance [61, 62]. While most systems that exhibit oscillatory interlayer coupling have in-plane magnetic anisotropy, in recent years multilayers
consisting of Co/Pd, Co/Pt structures have attracted attention due to their perpendicular magnetic anisotropy and thus the potential for ultrahigh density data recording applications [63-66]. For instance, investigations have shown that the Co/Pt multilayer structures with perpendicular anisotropy exhibit Ruderman-Kittel-Kasuya-Yosida (RKKY) oscillatory interlayer coupling with a ferromagnetic background between adjacent Co layer as the NM layer thickness increases [67, 68]. The overall magnetization reversal properties of such multilayer structures would thus be influenced by the interlayer coupling and the response of the individual Co layers to the reversing field.

In this study, microscopic Kerr imaging was used to study the perpendicular magnetization reversal of Co/Pt multilayers and their dependence on the strength and nature of the interlayer coupling between neighboring Co layers. Unlike conventional magnetometry such as superconducting quantum interference device (SQUID) and vibrating sample magnetometer (VSM) that reveal the overall magnetic response of the structure to an external magnetic field, Kerr imaging has, as illustrated in this study, the advantage of yielding the response of individual Co layers with high spatial resolution.

6.2 Methodology

Several series of \([\text{Co}(4 \text{ Å})/\text{Pt}(t_{\text{Pt}})]_N\) multilayers with repetition \(N\) from 5 to 30 were fabricated using a ultrahigh vacuum magnetron sputtering system as described elsewhere [67]. 50 mm long Si wafers with a native oxide layer were used as substrates. A 100 Å Pt buffer layer was first deposited on each wafer, followed by the deposition of Co/Pt multilayers. Each Co layer is uniform with a thickness of 4 Å and each Pt layer is a wedge with thickness \(t_{\text{Pt}}\) ranging from 0 to 80 Å. Finally, a 30 Å Pt layer was
deposited on top as the capping layer. Each Si substrate was cut into 40 pieces of 1.25 mm wide strips. Each strip has a thickness variation of 2 Å and \( t_{Pt} \) refers to the average thickness of the strip. Hysteresis loops were measured with a magnetic field \((H)\) perpendicular to the film plane using a LakeShore VSM. As evidenced by hysteresis loops in a previous report, perpendicular anisotropy was observed in all \([\text{Co}(4 \ \text{Å})/\text{Pt}(t_{Pt})]_N\) multilayers \([67]\). In this chapter, Co/Pt Multilayers with \(N = 5, 8, \text{ and } 12\), and \(11 \ \text{Å} \leq t_{Pt} \leq 41 \ \text{Å}\) were studied using magneto-optical Kerr imaging to reveal the details of magnetic domain wall formation and propagation during magnetic switching processes.

In situ magneto-optical images based on the polar Kerr effect were recorded at room temperature using 488 nm laser illumination \([69]\), polarizing optics and magnetic coils capable of producing magnetic fields perpendicular to the sample surface. The evolution of domain patterns and the reversal of magnetization in the multilayers were mapped through the following sequence: A magnetic field far exceeding the coercive field was initially applied for 3 seconds to magnetically saturate the sample following which the field was reduced to zero. The resulting magnetic state - the initial configuration - appears as a dark Kerr image with “down” magnetization in Figs. 6.1c—6.3c. The sample is then subject to a succession of pulsed magnetic fields \((H_p)\) close to the coercive value and oriented opposite to the direction of saturation magnetization of the initial configuration, as indicated in Fig. 6.1b – 6.4b. This reversal field was created by a 30 ms square current pulse flowing through the magnetic coils. A progression of Kerr images of the magnetic domains was sequentially recorded between
all of the repetitive field pulses. The recording time for each Kerr image was 3 seconds in order to obtain good signal to noise quality. Eventually, a large magnetic field was applied to fully reverse the magnetization.

6.3 Experimental results

Fig. 6.1a shows the room temperature coercivity, \( H_C \), for the \( N = 5 \) multilayer structure \([\text{Co}(4 \text{ Å})/\text{Pt}(t_{\text{Pt}})]_5\) as a function of \( t_{\text{Pt}} \) lying between 3 and 60 Å. A clear peak in \( H_C \) is evident at \( t_{\text{Pt}} \sim 23 \text{ Å} \), and the oscillatory behavior in \( H_C \) has previously been associated with the RKKY coupling between Co layers through the Pt [67, 68]. The arrow in figure 6.1a identifies the Pt layer thickness (11 Å) for the particular sample where the reversal was studied using Kerr imaging. The corresponding hysteresis loop of the sample is shown in Fig. 6.1b where the arrow indicates the magnitude of the pulse field \( H_p \) for the magnetic reversal studies. The interlayer coupling which involves both the exchange interaction and dipole interaction between Co layers is always ferromagnetic for the whole range of Pt thickness. While at \( t_{\text{Pt}} = 11 \text{ Å} \), the coupling strength is strong, the interlayer coupling becomes weaker at larger \( t_{\text{Pt}} \). Figs 6.1c – 6.1g provide some images of the magnetic domain patterns in time order during the application of the series of repetitive pulse field \( H_p \). Eventually, a large magnetic field was applied opposite to the initial magnetization to saturate the whole sample with “up”: magnetization as shown in Fig. 6.1h. A clear evolution from the fully saturated (dark) state with a single domain of “down” magnetization (Fig. 6.1c) to the fully reversed (bright) state of a single “up” domain (Fig. 6.1h) is observed. One observes in figures 6.1d to g the nucleation of the reversed magnetic domains (bright) with “up”
magnetization that grow steadily at the expense of the “down” domain by displacing the
domain boundaries and nucleating new reversed cells. While a domain wall between an
“up” and a “down” domain can move - thereby allowing one domain to grow and another
to shrink, defects and inhomogeneities within the multilayers can serve as pinning sites
which freeze the domain wall. The domain wall will eventually be unpinned when the
net gain of Zeeman energy due to the growth of the opposite domains overcomes the
domain wall pinning energy. The same descriptions apply to Fig. 6.1h to m showing Kerr
images on the same sample and same location with magnetization reversing in the
opposite direction: From fully “up” to fully “down” using the same pulsing technique
mentioned above. Nucleation sites for the two opposite magnetization reversal directions
match each other (Fig. 6.1d and i). We note that the images in figure 6.1c to m show only
two distinct contrasts – bright and dark – i.e. the “down” magnetization and the “up”
magnetization. Given that the total thickness (105 Å) of the [Co(4 Å)/Pt(11 Å)]₅
multilayer plus Pt cap layer is smaller than the skin depth of the probe laser beam
(estimated to be ~124 Å), it follows that in this case all five Co layers within a given
spatial area switch at the same time. As the reversed magnetization is nucleated at a
specific region of a given layer, the strong interlayer coupling ensures that the layers
above and below within the same region also reverse. The same behavior of the magnetic
switching was observed for the same series of multilayers (N =5) with tₚt = 21 Å and 31
Å (Fig. 6.5 and 6.6), where the Co layers are still strongly coupled and the structure
reverses as a single ferromagnetic unit.
Figure 6.1: (a) Room temperature coercivity ($H_C$) for $[\text{Co}(4 \text{ Å})/\text{Pt}(t_p)]_N$ multilayers with repetition $N = 5$. (b) Hysteresis loop for sample with $N = 5$, $t_p = 11$ Å (arrow in (a)). (c) – (g): Some Kerr microscopy images of $[\text{Co}(4 \text{ Å})/\text{Pt}(11 \text{ Å})]_2$ recorded at consecutive number of field pulses at $H_p$ indicated in (b) created by current pulse flowing through field coils. (i) to (l): Kerr images with pulses $H_p$ applied in the reversed direction. (h) and (m): Kerr images of the multilayer after saturated by a large magnetic field opposite to the initial magnetization.
This observation of reversal as a single ferromagnetic unit is in contrast to the
domain patterns during magnetization reversal from a multilayer sample with the same
number of Co layers \((N = 5)\), but a greater separation \((t_{Pt})\), and thus weaker coupling
between adjacent Co layers. Figure 6.2 shows the evolution of the domain patterns from
the structure \([\text{Co}(4 \text{ Å})/\text{Pt}(t_{Pt})]_5\) for \(t_{Pt} = 41 \text{ Å}\). As in Fig. 6.1, ferromagnetic interlayer
coupling is corroborated from Fig. 6.2a. Due to the greater value of \(t_{Pt}\), a weaker
interlayer coupling than that in the \(t_{Pt} = 11 \text{ Å}\) sample (Fig. 6.1b) is suggested from Fig.
6.2b with a smaller \(H_C\). In these weakly coupled layers the Kerr images show that the
entire structure does not switch as a single ferromagnet but rather the Co layers reverse
independently. As more field pulses were applied at \(H_p\) indicated in Fig. 6.2b, four
distinct domains (labeled 1, 2, 3, 4) with clearly distinguishable colors are observed (Fig.
6.2e) in this multilayer. Further increases in the field ultimately aligns all magnetic
domains and there is no contrast between them (Fig. 6.2h). Although the hysteresis loop
in Fig. 6.2b shows a sharp magnetic switching which suggests that all the Co layers
reverse together, Kerr images reveal that the Co layers switch independently. The
magnetization of a region within a Co layer with the lowest coercivity begins to reverse
first (bright regions Fig. 6.2e) and forms a “up” domain. This is followed by a layer-by-
layer reversal as evident by the shades of distinct different gray in Fig. 6.2d – 6.2g when
more field pulses causes regions with a little higher coercivity to reverse. Ultimately,
except for a few defective inclusions on the surface, the fully bright region develops in
Fig. 6.2h when positive saturation of the entire structure is achieved. The same
descriptions apply to Fig. 6.2h to m showing Kerr images on the same sample and same
location with magnetization reversing in the opposite direction: From fully “up” to
fully “down” using the same pulsing technique mentioned above. Four distinct domain colors (1, 2, 3, 4) are observed in Fig. 6.2k. Note that nucleation sites for the two opposite magnetization reversal directions do not match each other (Fig. 6.2e and k). Explanation to this is given later in this section.
Figure 6.2: (a) Room temperature coercivity ($H_C$) for $[\text{Co}(4 \, \text{Å})/\text{Pt}(t_{\text{Pt}})]_N$ multilayers with repetition $N = 5$. (b) Hysteresis loops for sample with $N = 5$, $t_{\text{Pt}} = 41$ Å (arrow in (a)). (c) – (g): Some Kerr microscopy images of $[\text{Co}(4 \, \text{Å})/\text{Pt}(41 \, \text{Å})]_5$ recorded at consecutive number of field pulses at $H_p$ indicated in (b). (i) to (l): Kerr images with pulses $H_p$ applied in the reversed direction. (h) and (m): Kerr images of the multilayer after saturated by a large magnetic field opposite to the initial magnetization.
Figures 6.3 show the response of similar, weakly coupled, Co layers \((t_{Pt} = 41 \text{ Å})\) with \(N = 8\). The respective thickness of the Co/Pt multilayers plus the Pt cap layers is 390 Å which is larger than the skin depth of the probing 488 nm laser radiation. Thus the rotation of polarization axis of the reflected light recorded in the Kerr images is influenced only by the magnetic state of the top several layers of the multilayers that lie within the penetration depth of the probe light. Once again, as the field pulses at \(H_p\) were applied, the lowest coercivity region begins to switch and grow while, depending on the local coercivity, regions of the other weakly coupled layers sequentially reverse quasi-independently. In Fig. 6.3e, four shades of distinct colors (1, 2, 3 and 4) are evident, revealing the four local magnetic states with different number of Co layers reversed. As more field pulses were applied, these domains are observed to grow as their magnetization aligns with the field. Eventually the stable state of the entire multilayer structure is reached when the magnetization of all domains reverses completely by and lies parallel to a large external magnetic field, leading to the uniform bright images with a single “up” domain (Fig. 6.3h). The same descriptions apply to Fig. 6.3h to m showing Kerr images on the same sample with magnetization reversing in the opposite direction: From fully “up” to fully “down” using the same pulsing technique mentioned above. Four distinct domain colors (1, 2, 3, 4) are observed in Fig. 6.3j.
Figure 6.3 (a) Room temperature coercivity ($H_C$) for [Co(4 Å)/Pt($t_{Pt}$)]$_N$ multilayers with repetition $N = 8$. (b) Hysteresis loops for sample with $N = 8$, $t_{Pt} = 41$ Å (arrow in (a)). (c) – (g): Some Kerr microscopy images of [Co(4 Å)/Pt(41 Å)]$_8$ recorded at consecutive number of field pulses at $H_p$ indicated in (b). (i) to (l): Kerr images with pulses $H_p$ applied in the reversed direction. (h) and (m): Kerr images of the multilayer after saturated by a large magnetic field opposite to the initial magnetization.
Fig. 6.4 shows a “bright” to “dark” series of images on Co/Pt sample with N=12 and $t_{\text{Pt}}=41\text{Å}$, with the most clear evidence of the independent switching of Co layers is shown in Fig. 6.4e – 6.4g where a darker shade appear when two domains with lighter shade overlap, analogous to putting two sunshades on each other to make a darker sunshade. The two overlapping domains are completely independent and grow individually. Within a region of the multilayer, the magnetization of one domain with some Co layer(s) reversed has no effect on the overlapping domain from other Co layer(s).
Figure 6.4: (a) Room temperature coercivity ($H_C$) for [Co(4 Å)/Pt($t_{Pt}$)]$_N$ multilayers with repetition $N = 12$. (b) Hysteresis loops for sample with $N = 12$, $t_{Pt} = 41$ Å (arrow in (a)). (c) – (h) Some Kerr microscopy images of [Co(4 Å)/Pt(41 Å)]$_{12}$ consecutive number of field pulses at $H_p$ indicated in (b). (i): Kerr images of the multilayer after saturated by a large magnetic field opposite to the initial magnetization.

It is evident that the magnetic orientation of the perpendicularly magnetized films has been determined with relatively high spatial resolution. With external field applied opposite to the magnetization of the fully magnetized multilayers, the formation of small domains of reversed magnetization is observed. We note that the hysteresis loops (Fig. 6.1b – 6.4b) of the four samples utilized in this study display abrupt switching transition with no evidence of the layer-by-layer magnetization reversal and do not reveal any peculiar signatures of the sequential reversal for the weakly coupled Co layers.
evident in the Kerr images in Fig. 6.2 – 6.4. The hysteresis loops alone cannot distinguish the difference of magnetic switching within the same series, e.g., $N = 5$, between $t_{Pt} = 11 \, \text{Å}$ (Fig. 6.1b) and $t_{Pt} = 41 \, \text{Å}$ (Fig. 6.2b). Thus it is evident that although a number of (meta) stable magnetic domain configurations are evident in the Kerr images for large $t_{Pt}$ that gives rise to weak interlayer coupling, the contribution of these states to the measured coercivity is not revealed in the hysteresis loops.

Figure 6.5 and 6.6 shows Kerr images of $[\text{Co}(4 \, \text{Å})/\text{Pt}(23 \, \text{Å})]_5$ and $[\text{Co}(4 \, \text{Å})/\text{Pt}(33 \, \text{Å})]_5$ multilayers correspondingly with magnetization reversal shown in both direction: from completely downward to completely upward (Fig. 6.5 a to f and Fig. 6.6 a to f) and from completely upward to completely downward (Fig. 6.5f to l and Fig. 6.6 f to l). These images are obtained by the same pulsing procedure as described before. Pattern of nucleation of reversed domain and their growth show no difference from the $[\text{Co}(4 \, \text{Å})/\text{Pt}(11 \, \text{Å})]_5$ (Fig. 6.1c to g), with only two distinct colors (dark and bright) identifiable, indicating that at these separations of Co layers (23 Å and 33 Å), the cobalt multilayers still reverse as a single ferromagnet. Interlayer coupling is not weak enough to allow layer by layer reversal to occur. It is also noted that nucleation sites for magnetization reversal in opposite direction match each other (Fig. 6.5b and g and Fig.6.6b and g).
Figure 6.5: Kerr images of magnetic reversal for [Co(4 Å)/Pt(23 Å)]$_5$ multilayers.
Figure 6.6: Kerr images of magnetic reversal for $[\text{Co}(4 \text{ Å})/\text{Pt}(33 \text{ Å})]_5$ multilayers.

Fig. 6.7 and 6.8 shows Kerr images of $[\text{Co}(4 \text{ Å})/\text{Pt}(41 \text{ Å})]_{12}$ for two cycles of round trip magnetization reversal, taken on another location of the same sample as in Fig. 6.4. Layer by layer magnetization reversal is observed as four distinct colors (four distinct state of magnetization), as labeled in Fig. 6.7d and j and in Fig. 6.8a and j. This behavior is consistent with our previous results (Fig. 6.2 to 6.4). Layer by layer reversal
are observed both in the dark to bright Kerr image series (Fig. 6.7a to f and Fig. 6.8a to f) and in the bright to dark image series (Fig. 6.7g to l and Fig. 6.8g to l). This indicates that the physics underlying the magnetic reversal is symmetric in the two opposite directions (downward to upward and upward to downward). However, in contrast to the sample of $[\text{Co}(4 \text{ Å})/\text{Pt}(23 \text{ Å})]_5$ and $[\text{Co}(4 \text{ Å})/\text{Pt}(33 \text{ Å})]_5$, nucleation sites during the reversal appear to be random. Comparing Fig. 6.7b, h, and Fig. 6.8a, g, there do not exist any matches in the nucleation location. Since cobalt layers are only weakly coupled, energy barriers for magnetization reversal are small and nucleation is more susceptible to intrinsic fluctuations during each individual process of magnetic reversal (e.g. exact shape of the magnetic field pulses) and leads to randomness in the nucleation site. We also see clearly in Fig. 6.7 and 6.8 that defects on sample surface act as obstacles for domain wall motion.
Figure 6.7: Kerr images of magnetic reversal for [Co(4 Å)/Pt(41 Å)]$_{12}$ multilayers.
Figure 6.8: Kerr images of magnetic reversal for [Co(4 Å)/Pt(4 Å)]_{12} multilayers.

6.4 Conclusion

In summary, a sequence of short pulses of externally applied magnetic fields are used to reverse the magnetization of [Co(4 Å)/Pt(t_{Pt})]_{N} multilayers with perpendicular magnetization. Monitored by magneto-optical Kerr imaging, the whole multilayer is observed to switch as a single ferromagnet for small values of Pt layer thickness when the Co layers exhibit strong coupling. As the Co layers become weakly coupled, clear
evolution of (quasi) independent switching of the individual layers is observed as more field pulses is applied. Magneto optical imaging reveals details of the layer-by-layer reversal through the appearance of magnetic domains with distinct shades colors for weakly coupled Co layers at large Pt thickness, enriching our understanding of the magnetic reversal mechanism of Co/Pt multilayers with oscillatory interlayer coupling.
CHAPTER 7

CONCLUSION

In this thesis research it has been demonstrated, in several relatively independent projects (Chapter 4, 5 and 6), the entire process of design, growth, fabrication and light scattering characterization of novel magnetic metallic micro-structures. While precious experience was acquired in growing and fabricating these structures, it was shown that light scattering techniques can be an exclusive tool in addressing fundamental physics issues related to magnetism in micron and sub-micron systems. For example, in Chapter 4, by exploiting Brillouin light scattering, current induced Amperian field modification of spin wave mode profile was revealed. The traditional method of propagating spin wave spectroscopy (PSWS) [70] to probe spin waves, involves electrical detection at microwave frequencies, a technique that is insensitive to these changes in oscillation profile. Ferromagnetic resonance, on the other hand, is also inadequate as it simultaneously probes the spin wave signal contribution from both magnetic layers in a bilayer structure. Only light scattering, because of its finite penetration depth, is capable of monitoring spin wave oscillation in individual magnetic layers.

Another example demonstrating the unique capability of light scattering techniques is the observation of layer by layer magnetic reversal of Co/Pt
multilayers presented in Chapter 6. Bulk detection methods, such as vibration sample magnetometry, display sharp magnetic transitions in hysteresis loop for both coupled and decoupled multilayers and reveal no information on the individual layer response during magnetic switching. Hysteresis loops from bulk magnetometry also reveal a gradual decrease of coercive field against increasing platinum thickness and do not reveal at which separation the cobalt layers are decoupled. On the other hand, the technique of Kerr imaging distinguishes the differences between coupled and decoupled multilayers during magnetic switching and identifies the specific separation of cobalt layers beyond which decoupling occurs. It is interesting to notice that Spin-polarized Scanning Electron Microscope (Spin-SEM) [71], which is a powerful tool in magnetic imaging, is not capable of observing such layer by layer reversal since the probing depth of spin-SEM is only 10Å, which is much smaller than the thickness of typical multilayers (more than 100Å).

In Chapter 5, it was demonstrated that the phenomenon of spin-polarized current induced domain wall movement can be observed by monitoring the polarization of scattered light (Kerr imaging). We have observed the motion of domain walls which are correlated with the direction of applied current through the channel. Occasional violations of correlation between wall movements and current flow directions are likely due to relatively wide permalloy channels (8µm) utilized and the nature of the domain wall. The 50nm channel film thickness leads to strong interference of domain wall movement from magnetostatic interaction inside the channel. It is suggested that future experiments should be repeated with thinner channels of 10 nm thick in order to suppress magnetostatic interactions. It is important that the surface quality of the SiO₂/Si
substrate should be improved to enable these thinner channels become suitable for the
domain wall movement experiments of interest. Synchrotron radiation facilities at
Argonne National laboratory [59] provide opportunities in probing magnetic wires with
sub-micron width, which are free from the interference of local magnetostatic interaction
and is likely to yield more conclusive results. Observation of spin-polarized current
induced domain wall movement with light scattering techniques is definitely possible and
study along this direction should be continued.

As evident from the results of chapters 4, 5 and 6 the advantage of light scattering
reside in their finite penetration depth in magnetic metallic materials and sensitivity to
high frequency oscillations in these materials. New directions of research in light
scattering include the use of femtosecond pulsed laser beams, which will provide ultrafast
time resolution [28-30]. The high frequency properties of these magnetic metallic
structures are important in achieving rapid magnetic switching since present computer
clock pulses are approaching frequencies in the same range. Near field scanning optical
microscope (NSOM) for magneto-optical measurement purpose was developed to
overcome the diffraction limit of light resolution [31, 32]. Another important trend is use
of light scattering as a tool in the studies of “spintronics” in semiconductors [72]. As
mentioned in Chapter 1, spintronics is subdivided into two main categories: magnetic
metallic substance and semiconductor systems. While magnetic metallic materials
provide memory applications, spintronics in semiconductors contributes largely to
computation. Injection of spin polarized electron into semiconductors produce spin
population imbalance and thereby to magneto-optical effects. In the case of light
scattering, important information on these spin-polarized electrons within the
semiconductor (oscillation frequency, lifetime) can be revealed by means of time resolved Kerr effect [4]. Such information will be important for constructing new generation of electron spin (instead of electron charge) based computation. Light scattering will remain a major complement to the mainstream, all-electrical, detection in the development and advancement of the physics of spintronics.
BIBLIOGRAPHY


We note the small discrepancy between the calculated SO and SA mode frequencies (12.3 and 15.3GHz) and the experimental observed values (12.0 and 14.5GHz). This difference probably arises from uncertainties in the exact thickness of the NiFe and Cu layers, plus deviation of the saturation magnetization from the nominal value of 812Oe, a quantity that is sensitive to the precise composition ratio of Ni to Fe. For example, we verified that modifying the Cu, NiFe thickness and the saturation magnetization by less than 6% will bring the calculated surface and volume mode frequencies into
agreement with the experimental values. This difference between the calculated and experimental values of surface mode frequencies does not affect any of the essential physics underlying this study.


[44] Note that absolute intensity values from different Brillouin spectra are not exactly comparable. The parameter S is defined to be insensitive to such variations. We impose the constraint $0.25 < R < 4.00$ to account for the intensity resolution in our BLS spectra.

[45] Since $\omega_h(\omega_l)$ monotonically increase(decrease) with $\Delta g$, the condition on the frequencies can be satisfied by merely tuning the magnitude of $\Delta g$.


[53] The number of pulses between each image is chosen to two to produce the greatest movement of domain wall between images while the time interval between each image is taken is minimized.


[56] Request for the LabView and Visual C++ program should be
sent to the e-mail address au.24@osu.edu.

[57] *The fast fourier transform and its applications*, written by E. O. Brigham, (Prentice-Hall, New Jersey, 1988)


