MASTER

Spin Hall effect induced magnetization dynamics in Pt/CoB/Pt nano-wires

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Abstract

Over the last years, current induced magnetization manipulation has become increasingly important in data storage applications. In the development of these new memory devices, there is an important role for perpendicular magnetized materials, especially in domain wall based applications. Current induced dynamics in those perpendicular magnetized materials are very complex. In recent current assisted de-pinning experiments, a current induced field like effect was observed. An effect with the same symmetry was measured in current assisted magnetic switching experiments. Two current induced effects arising from spin-orbit coupling and capable of exerting a torque that could explain those observations are proposed: the spin Hall effect (SHE) and Rashba spin-orbit coupling. This research focused on Pt/CoB/Pt systems for which the Rashba effect can be neglected, which allows us to focus on the SHE. We performed two types of measurements to characterize the SHE in our systems.

We started with current assisted field induced switching experiments. We used macro-spin simulations to show that the SHE is capable of assisting field induced magnetic switching. We fabricated Pt/CoB/Pt nano-wires and measured the switching field as a function of the applied current. Those measurements confirmed that we are able to assist or oppose the field induced magnetic switching, depending on the current and in-plane magnetic field direction, as was predicted by the simulations. This confirmed the presence of the SHE in our systems. We also tried to switch the magnetization using only the spin Hall effect in combination with an applied in-plane magnetic field. We clearly observed the current induced effect with the correct symmetry, but we were not able to systematically study the SHE using those measurements, because the in-plane field brings the system in a poorly defined multi-domain state.

Subsequently, we investigated a new technique to study the SHE. Using the results we obtained from the switching experiments, we performed macro-spin simulations that predict an influence of the SHE on the magnetic damping. This means that we could perform damping measurements using a TR-MOKE set-up to study the SHE. This experiment has many critical points, and the maximal current we can send through the nano-wire is limited by Joule heating. We put a lot of effort in optimizing the signal, but it was not possible to obtain a good enough signal to noise ratio to measure a current induced variation of the damping.
## Contents

### 1 Introduction
1.1 Metal ferromagnetism in thin films and recording media .................................. 1  
1.2 Spintronics in memory devices ................................................................. 3  
1.3 This thesis .................................................................................................. 7  

### 2 Theory
2.1 Magnetic energy and the effective magnetic field ........................................ 14  
2.1.1 Magnetic energy .................................................................................. 14  
2.1.2 Magnetic anisotropy ......................................................................... 15  
2.2 Magnetization dynamics .......................................................................... 21  
2.2.1 The LL and the LLG equation ............................................................. 21  
2.2.2 Damping ........................................................................................... 22  
2.3 Current influenced magnetization dynamics ........................................ 26  
2.3.1 Spin transfer torque ......................................................................... 27  
2.3.2 Spin Hall effect ................................................................................ 32  
2.3.3 Rashba spin-orbit coupling ................................................................. 36  
2.4 Simulations of current induced magnetization dynamics ................... 37  
2.4.1 SHE induced magnetic switching simulations .................................. 38  
2.4.2 SHE influenced magnetic damping simulations ............................. 40  

### 3 Methods
3.1 Sample fabrication and design ................................................................. 43  
3.1.1 Magnetron sputtering for thin film deposition .................................... 43  
3.1.2 Sample patterning ............................................................................. 44  
3.1.3 Oxidation ............................................................................................ 46  
3.1.4 Sample design .................................................................................... 47  
3.2 Measurement methods ............................................................................. 50  
3.2.1 Magneto optical Kerr effect ............................................................... 50  
3.2.2 Ultrafast demagnetization ................................................................. 52  
3.2.3 TR-MOKE Set-up .............................................................................. 53  
3.2.4 Kerr-Microscopy ............................................................................... 56
# 4 Current assisted switching measurements

4.1 Field-induced current-assisted switching measurements .......................... 60
  4.1.1 Decomposing the change in switching field ................................. 61
  4.1.2 Pt(5)/CoB(0.8)/AlOx(0.89) ........................................ 63
  4.1.3 Pt(5)/CoB(0.7)/Pt(2) wires .......................................... 64
  4.1.4 Pt(5)/CoB(0.7)/Pt(5) .................................................. 65
  4.1.5 Estimate of the effective Hall angle ..................................... 66

4.2 Current induced magnetic switching ................................................ 67
  4.2.1 Multi-domain state ......................................................... 68
  4.2.2 Current induced switching measurements .................................. 68

4.3 Contributions to the symmetric part of the switching field ................. 72
  4.3.1 Joule heating .............................................................. 72
  4.3.2 Oersted fields ............................................................. 73

# 5 Current influenced magnetic damping

5.1 Field depended damping measurements ............................................. 78
  5.2 Optimizing the pump beam power ............................................. 80
  5.3 Damping measurements on magnetic thin films ................................ 81
  5.4 Damping measurements with current induced spin torques ............... 83

# 6 Conclusions and outlook

6.1 Conclusions ................................................................. 89
  6.2 Outlook and applications ................................................ 90
Chapter 1

Introduction

Magnetism has interested and helped mankind for already thousands of years. From the first lodestone magnets invented by the Chinese, until the high tech memory devices in which magnetic materials are used today. Those magnetic memory devices developed with an incredible speed, starting from the magnetic core memories in the 50s until the modern MRAM. A mayor breakthrough took place in 1988 with the discovery of giant magneto-resistance (GMR). Together with some other observations, this was the beginning of the field of spintronics. In spintronics, not only the electric properties of the electron are used but also its spin. The magnetic moment associated with the spin can interact with the magnetization. Therefore it is possible to use an electric current to manipulate the magnetization of a magnetic material. In this first chapter, we will give an introduction to spintronics, and in particular to current influenced magnetization dynamics, which is the subject of this thesis. We will start with an introduction to ferromagnetism in recording media. In section 1.2 we will introduce the use of spintronics in magnetic memory devices. Subsequently, current assisted depinning experiments performed in this group will be discussed. Those measurements were the starting point of this project. In section 1.4 the research covered in this thesis and an outline of this report are discussed.

1.1 Metal ferromagnetism in thin films and recording media

Metal ferromagnets have a very important role in modern recording media. Before we can introduce those applications, we give a review of some important properties of ferromagnets which are fundamental for these applications.

We start with the discussion of the origin of magnetic properties in 3d ferromagnets. The valence band of the ferromagnetic metals iron (Fe), cobalt (Co) and nickel (Ni) contains delocalized 4s and more localized 3d electrons [1]. The wave-functions of
those electrons overlap, and therefore their atomic energy levels converge into energy bands. Due to the exchange coupling between spins, spins prefer to be aligned and this leads to a splitting of the spin-up and spin-down energy bands. In figure 1.1a the s and d band for Co are depicted. The states are filled till the Fermi level, and we see that there will be an imbalance in spin-up and spin-down electrons, which will lead to a net magnetic moment. The spins that are dominant are called majority spins, the other minority spins. By definition the magnetic moment points anti-parallel to the majority spin direction.

In models, the exchange coupling between spins responsible for ferromagnetism is threaded as an effective field, which we call the exchange field. However it is not a real field and it does not exert a torque on either the orbital moment of the electrons nor the nuclear magnetic moment. For this reason, the exchange coupling does not couple the spin system to the lattice, and therefore does not introduce a preferential direction for the magnetization. However, in most magnetic systems there is a preferred direction observed for the magnetization, called magnetic anisotropy. The preferred direction of the magnetization is called the easy axis. The axis along which the energy is maximized is called the hard axis. There are two contributions to this anisotropy: Shape anisotropy, and magneto-crystalline anisotropy. In chapter 2, we will come back to this important property of magnetic systems.

For magnetic thin films, the easy-axis can either lie in the plane of the sample or perpendicular to it. The former is called in-plane anisotropy, the latter is called out-of-plane anisotropy. Out-of-plane anisotropy is caused by an interface effect and is therefore observed in very thin films.

For magnetic memory, another important property of ferromagnets is the possibility to break in a multi-domain state. Sometimes the system can minimize its magneto-static energy by breaking-up into magnetic domains, see figure 1.1b. The domains are separated by domain walls. In the domain wall, the magnetization turns from one direction into the other direction. These domains cost a certain amount of energy to create, for two reasons: neighboring spins in the domain-wall are not aligned, so the exchange energy is not minimized. This term wants to create extended domain walls to minimize the exchange energy. The second contribution is from the anisotropy: the magnetization in the domain wall is not aligned with the easy axis. This term wants to create narrow domain walls. Because the domain walls costs energy to create, there will be a competition between minimizing magneto-static energy by creating domains, and minimizing domain-wall energy by minimizing the amount of domains.

Those domains play a crucial role in magnetic memory: by dividing a magnetic film into domains, we can assign a digital bit to the magnetization direction of a certain domain. In traditional hard disk drives (HDDs), magnetic films with in-plane magnetic anisotropy were used to store data. Over the years, people kept decreasing the size of the magnetic domains to increase the bit density of the memory devices. However, at some point the magnetic domains become unstable: already at room
1.2 Spintronics in memory devices

In the previous section, we introduced magnetic domains, which can be used in memory devices to store information. In early HDDs the information was read out by the induced current in a coil when the magnetic domains pass the read-head. Later, the anisotropic magneto-resistance of magnetic materials was used. This property allows for determining the state of a magnetic multilayer by measuring its resistance. In the late 80’s, a major breakthrough took place with the discovery of the Giant Magneto Resistance (GMR), which not only lead to drastically improved read-heads in HDDs, but gave birth to a whole new fascinating research field: spintronics. In this section we will give an introduction to the field of spintronics, starting with an explanation of GMR and its application in magnetic read-heads. Also another important discovery in spintronics is introduced: spin-transfer-torque (STT). STT has opened a whole new
world of opportunities for magnetic memory devices.

The beginning of spintronics: GMR

In 1856, Lord Kelvin discovered what is now known as the anisotropic magneto-resistance effect (AME) [4]. He discovered that the electrical resistance of iron increases when he applied a field collinear to the current, and decreases with a field applied perpendicular to it. This effect was gratefully used in magnetic read-heads until the 90’s: it allowed for the detection of the bits in a magnetic hard-drive. In the late 1980’s, Fert and Grunberg [3] independently discovered a much larger kind of magneto-resistance, later called the Giant Magneto Resistance (GMR). To explain this effect, we need to consider not only the charge properties of the electron, but also its spin angular momentum, which represents a certain amount of magnetic moment. Spintronics was born.

To explain the GMR effect, we start with a short explanation of electron transport in ferromagnetic metals [5]. Electrical current in a metal is mainly carried by the itertant s electrons, because they have a much lower effective mass than the d electrons. When the electrons are transported through the material, they can scatter, and the scattering probability determines the resistance. Because ferromagnets have a partially filled d shell there is a high number of available states at the Fermi level of the d band to which the s electrons can scatter. Because the probability of these transitions is much larger than for ordinary scattering, this leads to an enhanced damping. Because the density of states is different for majority and minority electrons, the scatter probability and thus the mean free path will be different for both types of electrons. This is illustrated in figure 1.2a.

Electrons that move through the material can also undergo spin-flip scattering, a process during which the electron will flip its spin. However, if the dimensions of the system are much smaller than the typical spin-flip scattering length, we can pretend that there are two independent spin channels in the material: one for spin-up, and one for spin-down electrons. This allows us to calculate the total resistance with a resister network replacement scheme, see figure 1.2b. For a single magnetic layer, it turns out that the resistance is independent of the magnetization direction of the sample.

However, when we add a second magnetic layer and change the relative orientation of the magnetization of those two layers from parallel to anti-parallel, we will find a change in total resistance. This is called the Giant Magneto Resistance (GMR). The change in resistance can be evaluated using the resistor scheme to be:

\[
\frac{\Delta R}{R} = \frac{(R_H - R_L)^2}{4R_H R_L},
\]

in which \(R_H\) is the resistance for a electron moving through a layer with its spin anti-parallel to the magnetization, and \(R_L\) the resistance for a electron moving through a
1.2. SPINTRONICS IN MEMORY DEVICES

Figure 1.2 – a) Cartoon of spin dependent electron transport in a ferromagnet. The large arrow indicates the direction of the magnetization of the material. The mean free path of the electrons depends on the orientation of the electron spin relative to the magnetization. The reason is the large amount of free states at the Fermi level for minority electrons at the Fermi level, to which the electrons can scatter. b) Resistor model of the GMR effect.

Layer with its spin parallel to the magnetization.

The above explained GMR is intensively used in read heads in nowadays memory devices. In figure 1.3, the working principle of the GMR read-head is explained. The simplest version of the GMR read-head contains two ferromagnetic magnetic layers and an anti-ferromagnetic layer. The anti-ferromagnetic layers pins the direction of the first ferromagnetic layer due to exchange coupling. This layer is therefore called the pinned layer. This layer is very hard to switch. The second ferromagnetic layer can be switched much easier. As discussed before, the resistance of the stack depends on the relative orientation of the magnetization of the two layers. When the read-head is brought in a magnetic field, the free layer will align with the field direction. This field could be the stray field from a magnetic sample, like in figure 1.3a. By measuring the resistance, we determine the orientation of the magnetization of the sample. In this way, we can read out the bits stored in magnetic domains by scanning the read head over the magnetic film, see figure 1.3b and c.

Current induced domain wall motion

As mentioned before, the discovery of GMR is often seen as the birth of spintronics. Another milestone in the development of spintronics was the discovery of spin transfer torque (STT) [7], which states that the spins in a spin polarized current can exert a
**Figure 1.3** – Basics of the GMR read-head a) Layout of a GMR read-head. The free layer will align its magnetization with the stray field of a magnetic bit. b) and c) By measuring the resistance of the stack we can read out information stored in magnetic bits.

**Figure 1.4** – Domain wall motion by spin transfer torque. The magnetic moment associated with the electron spin is not parallel to the magnetization in the domain wall, and thus exerts a torque on the magnetization. In this way the domain wall will be moved along with the electron.
torque on the magnetization when they are not parallel to this magnetization. In this way we can move magnetic domain walls through a magnetic wire with a spin polarized current, as sketched in figure 1.4. This idea inspired Parking [2] who envisioned a memory device which combines the low costs of HDDs and the high speed of solid state memory. Instead of moving the disc containing the magnetic bits mechanically to the read/write-heads in traditional HDDs, he proposed to move the magnetic bits through the material using STT. This would be both faster and more reliable (the disk in a HDD spins at high speed just a few nanometers below the write/read head). In figure 1.5a, a sketch shows the idea of this memory device, called the magnetic racetrack memory.

In the design of magnetic domain based memory devised like the racetrack memory, it is very important to be able to control the domain wall motion. Domain walls

### Magnetic domain wall de-pinning

In the design of magnetic domain wall based memory devises like the above introduced racetrack memory, it is very important to be able to accurately control the domain wall motion. The driving force behind domain wall motion can be an applied external field or for example the above introduced current induced STT. In figure 1.5b the domain wall velocity as a function of applied field is plotted. In a real crystal, the domain wall faces impurities in the crystal which act as pinning sites. When the driving force is low, the domain wall will get pinned at those pinning sites. When the temperature is larger then zero, thermal fluctuation can de-pin the domain wall. In this way, the domain wall moves from pinning site to pinning site. This is called the creep regime. This creep regime plays a crucial role in perpendicular magnetized materials, because it is very difficult to achieve domain wall motion beyond this creep regime.

A second reason that domain wall de-pinning is important, is the fact that we have to be able to very precisely control the position of magnetic domain walls in memory devices. Therefore, artificial pinning sites are created, which allow to control the position of domain walls. This means that we have to de-pin the domain wall to be able to move it. For this reason, current assisted de-pinning experiments are performed in this group.

### 1.3 This thesis

This last section introduces the work discussed in this thesis. In the first subsection, we will discuss previous work done in this group on current induced magnetic domain wall depinning which formed the starting point of this research. In the second subsection we will introduce the work discussed in this thesis.
Previous work

In the previous section the relevance of domain wall de-pinning for applications in memory devices was discussed. Also current induced STT was introduced, which could be used as the driving force behind domain wall de-pinning. In this group, current assisted field induced de-pinning experiments are performed to study the effect of this STT on the domain wall de-pinning. Those experiments were done on Pt/Co/Pt nanowires. On a small portion of the wire the anisotropy was lowered by bombarding the wire with ions [43]. This allows for the creation of domains, by first apply a large field in one direction, and then a smaller field in the other direction, which switches only the region with lowered anisotropy. Current pulses are now send trough the wire, and the de-pinning field is measured, using an optical set-up Figure 1.6 depicts this measurement method. The magnetization direction next to a domain wall is measured, and when a change in magnetization is observed, we know that the domain wall is de-pinned. Figure 1.7 depicts the expected and observed behavior of those current assisted de-pinning experiments. In the top drawing we see the initial situation with the three domains. If we apply a magnetic field that is larger than the de-pinning field in the direction of the domain in the middle, both domain walls get de-pinned and the domain starts to grow in both directions. When we send electrons through the wire from right to left, we expect due to STT that the current tends to drag the domain walls to the left. We assume that the current is low enough to be sure that the STT on itself is not strong enough to de-pin the domain walls. When we apply a field, we expect the left domain wall to be de-pinned at a field lower than the original de-pinning field,
1.3. THIS THESIS

Figure 1.6 – Using the TR-MOKE set-up to perform domain-wall de-pinning experiments. In a) the starting situation is given. When we apply an increasing magnetic field, at some point, called the de-pinning field $H_{dep}$, the domain wall will get de-pinned. If we measure the magnetization just next to the domain wall as depicted in b), we will be able to detect this de-pinning from the measured magnetic signal as depicted in c).

Figure 1.7 – Schematic representation of the expected and observed field induced current assisted de-pinning behavior. STT would drag both domains in the same direction, thereby helping the de-pinning of one domain and opposing the de-pinning of the other domain. The observed effect helps the de-pinning in both directions, and is therefore field like. The horizontal arrows indicate the direction of the electrons.
and the right domain should be de-pinned at higher fields, as depicted in the third image. However, a field like effect as depicted in the bottom image is observed. The current, depending on its direction, helps or opposes the de-pinning of domain walls in both direction (note that in figure 1.7 we only show the current direction for which we help the de-pinning. When the current direction is inverted, the current opposes the de-pinning of both domain walls). We call this effect field like, because it wants to stabilize one magnetization direction. It is not a real magnetic field. This observed behavior was very unexpected, and could not be explained by conventional STT.

Around the same time those experiments were performed, other groups also found current induced effects in comparable systems which had the same unexpected symmetry [16]. Two different explanations originating from spin-orbit (s-o) coupling were proposed to explain those observations: Rashba fields which are induced at the interface between the ferromagnet and the nonmagnetic metal, and the spin Hall effect (SHE), which creates a spin polarized current from the Pt layer into the magnetic layer. Both effects would exert a torque on the magnetization with the right symmetry to explain the measurements. More details about both effects will be given in chapter 2. In this thesis we are interested in the effect of those s-o torques on the magnetization dynamics in Pt/CoB/Pt systems. In the next subsection, we will introduce the work done during this research.

**This thesis**

In this thesis we want to investigate the current induced spin-orbit torques in multilayer nano-wires containing a magnetic layer and at least one Pt layer, which has large spin-orbit coupling. As will be discussed in chapter 2, in Pt(x)/CoB(x)/Pt(x) systems, where x indicates the thickness of the individual layers in nm, we can neglect Rashba fields, and therefore the main focus of this thesis is on the SHE. We performed two types of measurements to verify whether the current-induced spin-orbit torques are present and significant in Pt/CoB/Pt systems. To perform those measurements, we use a Time Resolved-MOKE (TR-MOKE) set-up which will be discussed in chapter 3. The first measurements, which will be discussed in chapter 4, were performed to investigate the effect of current induced s-o torques on magnetic switching. Those measurements are performed to investigate whether the predicted current induced s-o torques are present in our systems and to characterize them.

The second method is a new and alternative way to study the effect of the spin-orbit torque on the magnetization. When we bring the magnetization of the sample out of equilibrium, the magnetization will start a precessional motion which will damp out because of scattering processes. Both the Rashba field and the SHE could exert torques with components parallel or anti-parallel to this damping torque. For example, Ando et al [8] showed the manipulation of the magnetic damping in in-plane magnetized NiFe films by the use of the SHE in platinum. We want to use the TR-MOKE set-up
to measure the trajectory of the magnetization of out-of-plane magnetized materials as a function of time to determine the damping parameter. In this way we are able to investigate current induced effects on the magnetic damping, and give the opportunity to study the current induced spin-orbit torques.

We will now give an outline of this thesis:

- In chapter 2 we will start with a theoretical background of the various subjects covered in this thesis. We will discuss magnetization dynamics, and especially magnetic damping. We will also take a closer look to the current induced torques, the already introduced Rashba field, SHE and the STT. We will argue that the SHE could induce a torque on the magnetization that could influence the damping. We will also describe simulations that show the effect of the SHE on field induced magnetic switching and magnetic damping.

- In chapter 3, we will discuss methods used to fabricate the samples that are used in this thesis. Also the TR-MOKE set-up is explained in more details and we will introduce a second measurement technique, the Kerr microscope.

- In chapter 4, results on current assisted switching measurements will be discussed. Those measurements are performed to prove the existence of s-o torques in our systems and characterize the effect in different samples. We also performed measurements to study pure current induced magnetic switching.

- In chapter 5 we performed damping measurements on both thin films as well as nano-wires. We measured thin films to characterize our samples and to investigate the effect of the Pt layer thickness on magnetic damping. The measurements on the nano-wires were performed to study the effect of current induced s-o torques on the magnetic damping.

- In chapter 6 we will summarize the most important results of this thesis and give an outlook for possible future research.
Chapter 2

Theory

In this chapter we treat the theoretical background of current influenced magnetization dynamics. We will discuss which current induced effects we expect in our systems and how they can influence the dynamics of the magnetization.

In the first section, we will introduce magnetic energy and the effective magnetic field. The equilibrium position of the magnetization is determined by minimizing the magnetic energy, and we will see that magnetic anisotropy leads to a preferred magnetization direction. Interface effects will turn out to be important to obtain perpendicular magnetized systems, in which we are interested in this thesis.

In the second section we will take a look at the dynamics of the magnetization when the magnetization is brought out of equilibrium. Those dynamics are described by differential equations like the Landau Lifshitz or the Landau Lifshitz Gilbert equation. We will also introduce the Gilbert damping parameter $\alpha$ which describes damping in the LLG equation. This parameter is very important in this thesis because we are interested in damping measurements.

In the third section, current induced magnetization dynamics are introduced. Here we take a look how the dynamics introduced in section 2.2 are influenced by electrical currents. We will introduce spin transfer torque (STT) which is exerted on the magnetization when a spin polarized current is injected into a ferromagnetic system. Subsequently we take a closer look to the two spin-orbit torques introduced in chapter 1, which are proposed to explain the observed current assisted switching and domain wall depinning experiments. We will start with the spin Hall effect (SHE) in section 2.3.2, and cover the Rashba effect in section 2.3.3.

In the last section of this chapter we treat macro-spin simulations of the LLG equation including a STT induced by the SHE. Those simulations will show that indeed the SHE can assist field induced switching. In this thesis, we wanted to study the current induced spin-orbit torques by measuring their effect on the magnetic damping. Therefore we also performed simulations that demonstrate the effect of the STT induced by the SHE on the magnetic damping.
2.1 Magnetic energy and the effective magnetic field

Bits in magnetic memory devices are stored by controlling the magnetization direction of a magnetic domain. In the introduction chapter we mentioned that the size of magnetic domains in storage devices could be scaled down below 10 nm when systems with out-of-plane magnetization are used. Whether the equilibrium direction of the magnetization is in-plane or out-of-plane is determined by minimizing the magnetic energy of the system. This magnetic energy contains a contribution from the external applied field and from the anisotropy. We will start this section by introducing an expression for the magnetic energy. In the second subsection we will discuss the origin of magnetic anisotropy.

2.1.1 Magnetic energy

The energy of a free magnetic moment $\vec{\mu}$ in an external magnetic field is called the Zeeman energy:

$$E_{\text{zee}} = -\mu_0 (\vec{\mu} \cdot \vec{H}_{ex}),$$

in which $\mu_0$ is the permeability of vacuum and $H_{ex}$ is the external applied magnetic field. For a free magnetic moment, the energy is minimized when the magnetic moment is aligned with the applied magnetic field. Inside a crystal, a magnetic moment of an atom can interact with neighboring atoms. As we will see later on, those interactions can lead to contribution to the magnetic energy which depends on the direction of the magnetic moments and which we call magnetic anisotropy. We define this energy as $E_{\text{ani}}(\vec{M})$. This means that the magnetic moments will have a preferred direction, even without an applied magnetic field.

When we study the magnetic behavior of crystals we will not study individual magnetic moments, but the magnetization, which is defined as the volume density of magnetic moments:

$$\vec{M} = \frac{1}{V} \sum \vec{\mu}_i,$$

in which $V$ is the volume of the system and the summation is taken over all individual magnetic moments $\vec{\mu}_i$ inside the system.

The total magnetic energy can be written as the sum of the Zeeman energy and the anisotropy energy:

$$E_{\text{mag}}(\vec{M}, \vec{H}) = E_{\text{ani}}(\vec{M}) + E_{\text{zee}}(\vec{M}, \vec{H}).$$

We saw that a free magnetic moment will tend to align with an applied magnetic field. For the magnetization inside a crystal, we can define an effective field, along which the magnetization wants to align itself to minimize the magnetic energy. This effective field is defined as:
2.1. MAGNETIC ENERGY AND THE EFFECTIVE MAGNETIC FIELD

\[ \vec{H}_{eff} = -\frac{1}{\mu_0} \frac{dE_{mag}(\vec{M}, \vec{H})}{d\vec{M}}. \] (2.4)

This effective field is important when we are going to discuss the dynamics of the magnetization in section 2.2. In the next subsection we will take a closer look at magnetic anisotropy.

2.1.2 Magnetic anisotropy

In the previous section we introduced the anisotropic contribution to the magnetic energy. An axis along which the energy is minimized is called the easy axis, and an axis along which the energy is maximized is called the hard axis. Without this magnetic anisotropy, spontaneous magnetization in 2d systems would not be possible, and hardly observable in 3d systems. The reason is that the exchange interaction, responsible for ferromagnetic ordering, is short ranged. Therefore, it costs very little energy for the magnetization to rotate over a distance called the magnetic coherence length. In this way the system could minimize its magneto-static energy by curling the magnetization, which would lead to zero net magnetization. In this subsection we will take a look at the origin of magnetic anisotropy. We want to find an expression for the contribution of anisotropy to the magnetic energy given by equation 2.3.

The anisotropy is defined as the energy \( K \) it costs to turn the magnetization from the easy axis to the hard axis. In this thesis we work with thin films, so we will focus on those systems. In figure 2.1 we introduce the system of coordinates we use in this thesis. The thin film lies in the x-y plane, and z is the out-of-plane direction. In magnetic nano-wires, we define x as the direction of the wire. We define \( \Theta \) as the angle between the magnetization and the z axis. The angle between the applied magnetic field and the z axis is indicated by \( \beta \).

In first order, the magnetic anisotropy energy (MAE) is given by:

\[ E_{ani} = K \sin(\Theta). \] (2.5)

We observe that for positive \( K \), the energy will be minimized for an out-of-plane magnetization, and for a negative \( K \) for an in-plane magnetization. The anisotropy constant \( K \) can be written as the sum of two contributions: \( K = K_{mca} + K_s \). The first term is the magneto-crystalline anisotropy (MCA), the second term is the shape anisotropy. We will explain those two terms in a moment. This means that the combination of \( K_s \) and \( K_{mca} \) determines whether the easy axis is in-plane or out-of plane.

It was shown by Draaisma et al. [30] that the anisotropy of a thin film can also phenomenological be written as the sum of a volume and an interface term, which is
Figure 2.1 – System of coordinates used throughout this thesis. The out of plane direction is always along the z axis and a positive current direction is defined to be in the positive x direction.

more convenient when working with thin films:

\[ K = K_{\text{vol}} + \frac{2K_{\text{sur}}}{t}, \]

(2.6)
in which \( t \) is the thickness of the magnetic layer and the factor 2 enters because the magnetic layer is enclosed by two surfaces. This expression of the anisotropy constant shows that there is a competition between the volume and surface anisotropy which determines the sign of \( K \) and thereby the direction of the easy axis. This means that the anisotropy will be out-of-plane when:

\[ K_{\text{sur}} < -\frac{tK_{\text{vol}}}{2}. \]

(2.7)

So whether we observe perpendicular anisotropy in a system depends on the relative size of the volume and the surface anisotropy, and the thickness of the magnetic layer. The volume term includes shape anisotropy and a weak contribution of the magneto-crystalline anisotropy. The size of both contributions to the volume anisotropy for Fe, Co is Ni are listed in table 2.1. Note that the shape anisotropy is listed for thin films. The surface anisotropy is mainly due to magneto-crystalline anisotropy which is strongly enhanced at the surface due to interfacial bonding effects. We will now take a closer look at the physics behind the shape and magneto-crystalline anisotropy, and the reason the latter is enhanced at the surface.

Shape anisotropy

Shape anisotropy is the result of dipole-dipole interaction. The magnetic moments in the material feel the magnetic fields originating from other magnetic moments. In
2.1. MAGNETIC ENERGY AND THE EFFECTIVE MAGNETIC FIELD

<table>
<thead>
<tr>
<th>Metal</th>
<th>Structure</th>
<th>$K_s [eV/atom]$</th>
<th>$K_u [eV/atom]$</th>
</tr>
</thead>
<tbody>
<tr>
<td>Fe</td>
<td>bcc</td>
<td>$-1.4 \times 10^{-4}$</td>
<td>$-4.0 \times 10^{-4}$</td>
</tr>
<tr>
<td>Co</td>
<td>hcp</td>
<td>$-9.3 \times 10^{-4}$</td>
<td>$-5.3 \times 10^{-4}$</td>
</tr>
<tr>
<td>Ni</td>
<td>fcc</td>
<td>$-1.2 \times 10^{-4}$</td>
<td>$-8.6 \times 10^{-4}$</td>
</tr>
</tbody>
</table>

Table 2.1 – Bulk anisotropy energies for Fe, Co and Ni at 4.2 K. Both the shape anisotropy $K_s$ for thin films and the magneto-crystalline anisotropies $K_u$ are listed. Data adopted from [25].

First order the magnetic moments can be regarded as magnetic dipoles, and the energy associated with dipole-dipole interaction between two dipoles is given by:

$$E_{dip-dip} = -\frac{\mu_0}{4\pi r^3} \left[ \vec{m}_1 \cdot \vec{m}_2 - 3(\vec{m}_1 \cdot \vec{r})(\vec{m}_2 \cdot \vec{r}) \right],$$

with $r$ the distance between the two dipoles $\vec{m}_1$ and $\vec{m}_2$ and $\vec{r}$ a unit vector pointing from $\vec{m}_1$ to $\vec{m}_2$. In figure 2.2a we show four possible alignments between two magnetic moments and the associated dipole-dipole energy in multiples of $E_{dip-dip}^0$ which we define as $E_{dip-dip}^0 = \frac{\mu_0 m^2}{4\pi r^3}$. In figure 2.2b we show how this interaction leads to anisotropy in a magnetic system. The total dipole-dipole energy is the summation over all dipole-dipole interactions in the system and for simplicity, we only take nearest neighbor interactions into account. Figure 2.2b shows a magnetic sample containing eight magnetic moments. In the top drawing, the magnetization points in-plane, in the bottom drawing it points out-of-plane. The total dipole-dipole energy of both configurations is given and we observe that the in-plane configuration has lower energy. In general, shape anisotropy wants to align the magnetization along the longest axis of the system.

From a macroscopic point of view the effect of the dipole-dipole interaction can be described by an effective field which is called the demagnetization field $\vec{H}_d$. The demagnetization field is generally defined as:

$$\vec{H}_d = -\overline{N}\vec{M},$$

in which $\overline{N}$ is the demagnetization tensor which depends on the geometry of the sample. For a normal thin film, this tensor is given by [31]:

$$\overline{N} = 4\pi \begin{bmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{bmatrix}.$$  (2.10)

This means that only when $H_d \neq 0$, the magnetization has a component along the $z$ axis, and it tends to push the magnetization in-plane and thus introduces an in-plane anisotropy.
CHAPTER 2. THEORY

With the use of the definition of the effective field (equation 2.4) and the anisotropy constant, the dipole-dipole energy for thin films becomes:

\[ E_d = K_s = -\frac{1}{2\mu_0}M^2, \quad (2.11) \]

in which \( K_s \) is the shape anisotropy which is the main contribution to the volume anisotropy term in equation 2.6.

**Magneto-crystalline anisotropy and PMA**

We already mentioned that a second contribution to the volume anisotropy is due to magneto-crystalline anisotropy (MCA) and that this contribution is small in comparison to shape anisotropy. However MCA can be strongly enhanced at interfaces, leading to strong surface anisotropy which could lead to perpendicular magnetic anisotropy, as we will see in this subsection.

MCA is defined as a preferred magnetization direction with respect to the crystalline axis. The main contribution to magneto-crystalline anisotropy is spin-orbit coupling as explained by Van Vleck in 1937. Spin-orbit coupling couples the spin angular momentum of the electron to its orbital angular momentum, which is coupled to the lattice of the crystal.

We will first introduce spin-orbit coupling. The conventional Hamiltonian describ-
2.1. MAGNETIC ENERGY AND THE EFFECTIVE MAGNETIC FIELD

The energy of an electron in a magnetic and electric field is given by:

\[ H = \frac{1}{2m_e}(\hat{p} - \vec{A}) + V, \]  

(2.12)

in which \( \hat{p} \) is the momentum operator, \( (\hat{p} = -i\hbar \nabla) \), \( \vec{A} \) is the vector potential and \( V \) the scalar potential, describing the magnetic and electric fields. This Hamiltonian does not take relativistic effects into account. The relativistic version of the Schrödinger equation is called the Dirac equation. In first order, in the Dirac equation the Hamiltonian can be written as:

\[ H = \frac{1}{2m_e}(\hat{p} - A)^2 + V - \frac{1}{4m_e^2c^2}\vec{s} \cdot (\hat{p} \times \nabla V), \]  

(2.13)

in which \( m_e \) is the electron mass, \( c \) is the speed of light and \( \vec{s} \) represents the spin-direction. This extra term accounts for relativistic effects. When we use the definition of the quantum mechanical orbital momentum \( \vec{l} \), \( \hbar \vec{l} = (\vec{r} \times \hat{p}) \), and the electrostatic potential of the nucleus, \( V = Ze/(4\pi\varepsilon_0r) \), in which \( Z \) is the atomic number of the nucleus, \( r \) is the distance of the electron to the nucleus and \( \varepsilon_0 \) the permittivity of vacuum, we can rewrite equation 2.13:

\[ H = \frac{1}{2m_e}(\hat{p} - A)^2 + V - \frac{e\hbar^2}{2m_e^2c^2r} \frac{dV}{dr} \cdot (\vec{s} \cdot \vec{l}). \]  

(2.14)

The third term on the right hand side of the equation introduces spin-orbit coupling. The spin-orbit coupling part of the Hamiltonian is often written as:

\[ H_{so} = \xi(r) \cdot (\vec{s} \cdot \vec{l}), \]  

(2.15)

where \( \xi(r) \) is material dependent and determines the strength of the spin-orbit coupling. We will now explain how spin-orbit coupling leads to magneto-crystalline anisotropy.

In Co, Ni and Fe electrons in the unfilled 3d shell are responsible for ferromagnetism. The ten available states in the 3d shell can be written as a linear combination of 5 eigenfunctions of the \( L_z \) operator, which combined with the spin degree of freedom gives the ten possible states. Those 5 eigenstates are labeled with a magnetic quantum number \( m_l \) which can take the values -2, -1, 0, 1 and 2 and form a complete set of eigenfunctions. Those eigenfunctions are time dependent and thus carry angular momentum as indicated by their magnetic quantum number \( m_l \).

In the absence of any external magnetic field or other perturbation, those five eigenfunctions are degenerate and therefore any linear combination of those five states is still a valid solution to the Schrodinger equation. Therefore it is possible to construct a set of five stationary solutions, which are depicted in figure 2.3a. Those five states
still form a complete set of eigenfunctions. Note that those states are stationary and therefore carry no angular momentum.

When an atom is placed in a crystal, it is surrounded by other atoms, see figure 2.3b. We can describe the interaction of the atoms with their neighbors using the Ligand field model. In this theory we describe one central atom, that only feels the electron clouds of the nearest neighbors which are called Ligands, as depicted in figure 2.3b. The interaction between the electrons with the ligands is due to the Coulomb force. Because the coulomb force depends on the overlap of the electron wave-function with the electron cloud of the ligand, the energy of the orbitals depends on their spatial layout. This means that in this case not all orbitals will have the same energy but that their energy levels are shifted. The size of this shift in energy depends on the bonding strength between atoms. For very weakly bonded atoms, the energy shift will be very small. The spatial layout is only constant in time for the stationary orbitals depicted in figure 2.3a, which are therefore the eigen-functions of the Hamiltonian in the Ligand field theory.

Until now, we did not take spin-orbit coupling into account. The system can lower its energy by gaining a finite orbital momentum which couples with the electron spin. We argued that in the crystal, the stationary states are the eigenfunctions of the Hamiltonian. To gain a finite amount of angular momentum, the stationary states will have to mix. However, this mixing of states costs energy, because of the gain in electrostatic energy from the coulomb interaction. This means that the amount of mixing between states depends on the bonding strength between atoms.

This mixing of states can explain the large values of the surface anisotropy at the boundary of a magnetic and a non-magnetic material, for example Co and Pt. When we take again a look at figure 2.3a and define the z axis perpendicular to the interface between the magnetic and non-magnetic layer, we have two in-plane and three out-of-plane orbitals. At the interface, the two in-plane orbitals of the magnetic material feel the presence of other magnetic atoms of the same species. This leads to a certain amount of mixing of the two in-plane orbitals, depending on the bonding strength. The three out-of-plane orbitals feel the atoms of the neighboring material. The out-of-plane orbitals can also mix depending on the bonding strength between the magnetic and non-magnetic atoms. The bonding strength between the magnetic and non-magnetic atoms can be different from the bonding strengths between the magnetic atoms, and therefore the amount of mixing of the out-of-plane orbitals can differ from the mixing of in-plane orbitals. When the bonding with the non-magnetic atoms is stronger than the bonding in-between magnetic atoms, the out-of-plane orbitals are less mixed compared to the in-plane orbitals. Because in-plane orbitals lead to out-of-plane orbital momentum, this leads to a preferred out-of-plane direction of the orbital momentum, and thus to PMA. For example for a Co-Pt interface, the ratio between bonding strengths between Co-Pt and Co-Co bonds is 1.6 [32], leading to strong out-of-plane anisotropy at the boundary of Co and Pt.
2.2 Magnetization dynamics

In the previous section we described that in equilibrium, the magnetization lies along an effective field which is a combination of an external field and the anisotropy field. In this section we describe the dynamics of the magnetization when it is brought out of equilibrium. We will see that in this case the magnetization will start a damped precessional motion around the effective field. A number of differential equations exist which describe this behavior. We will discuss two of them in the first part of this section: the Landau-Lifshitz (LL) equation and the Landau-Lifshitz-Gilbert (LLG) equation. An important aspect of the LL(G) equation is the damping parameter. We will explain the origin of this parameter in the second part of this section.

2.2.1 The LL and the LLG equation

The torque on an isolated magnetic moment placed in an external magnetic field is defined as:

$$\vec{T} = \vec{\mu} \times \vec{H}_{ext},$$

(2.16)

where $\vec{\mu}$ is the magnetic moment and $\vec{H}_{ext}$ is the applied magnetic field. The effect of this torque on the magnetic moment is a precessional motion around the direction of the magnetic field. Inside a ferromagnet, the magnetization, which is defined as the total magnetic moment density, feels a similar torque. The difference being that the external field $\vec{H}_{ext}$ is replaced by an effective field $\vec{H}_{\text{eff}}$, introduced in the previous section. Using this effective field, Landau and Lifshitz [44] introduced an equation which describes the time evolution of the magnetization:
in which $\gamma$ is the gyromagnetic ratio defined as $\gamma = \frac{e g}{2 m_e}$, with $g$ the Landé factor, $e$ the elementary charge, $m_e$ the electron mass and $\mu_0$ the permeability of vacuum. This equation describes a precessional motion of the magnetization around the effective field direction. However, this formula fails to explain the damping process observed experimentally. Therefore, Landau and Lifschitz added a term proportional to $M \times [M \times H_{eff}]$ to account for damping. They defined a phenomenological damping parameter $\lambda$, which we will call the Landau-Lifshitz (LL) damping parameter. The LL-equation now becomes:

$$\frac{d\vec{M}}{dt} = -\gamma \vec{M} \times \vec{H}_{eff} + \lambda \vec{M} \times \vec{H}_{eff}. \tag{2.18}$$

Later, Gilbert [45] stated that the LL equation fails for large values of the damping parameter, and proposed to replace the LL damping term with a viscous force term:

$$\frac{d\vec{M}}{dt} = -\gamma_G \mu_0 \vec{M} \times \vec{H} - \frac{\alpha}{M_{sat}} \vec{M} \times \frac{d\vec{M}}{dt}, \tag{2.19}$$

with $\alpha$ the Gilbert damping parameter and $\gamma_G = \gamma (1 + \alpha^2)$. This equation describes the damping in ferromagnetic materials more accurately when the damping is large, like in thin magnetic films. With $\alpha = \lambda / \gamma$ both equations are mathematically equivalent. Until today, no complete agreement exists about which of the two equations describes best the damping in ferromagnetic materials [21]. In this report, we will assume that the LLG equation and the LL equation are equivalent and the damping parameters are related via $\alpha = \lambda / \gamma$. We have to make one more remark about the damping parameter. We will assume that the damping parameter is a scalar, although ab initio studies show that in almost all cases, the scalar damping parameter should be replaced by a tensor [20], because the damping parameter turns out to be anisotropic. As the LLG equation with the scalar form of the damping parameter describes very well the measurements done in this research and other studies in this group, we stick to the scalar form.

### 2.2.2 Damping

In the previous subsection we saw that the damping of the magnetic precessional motion towards its equilibrium position is described in the LLG equation by a parameter $\alpha$ called the Gilbert damping parameter. This parameter is linked to the damping parameter $\lambda$ used in the LL equation. In this subsection we will take a closer look at the mechanisms responsible for this damping. This subsection is split into two parts.
2.2. MAGNETIZATION DYNAMICS

2.2.2. Intrinsic mechanism

The intrinsic damping mechanism in metallic ferromagnets originates from a combination of spin-orbit coupling and electron-lattice scattering. Many theories exist which describe the intrinsic damping. One very intuitive model is the so-called breathing Fermi surface model[18]. We will start this section by explaining this model. As we will see, this model cannot explain the temperature behavior of the damping parameter around room temperature and higher for Co and Fe. Therefore we will introduce a second model, called the torque-correlation model, which is able to describe the temperature dependence of the damping in the whole temperature range. This model is however less intuitive and will only be discussed briefly.

The breathing Fermi surface model

As already mentioned, the mechanism responsible for intrinsic magnetic damping is spin-orbit coupling. When the magnetization precesses around the effective field, the energy contribution of the spin-orbit coupling changes and thereby the energy levels of the electronic states can shift. This means that some occupied states lying just below the Fermi level can be pushed above the Fermi level, and some empty states just above the Fermi level can be pushed just below the Fermi level (referred to as the breathing
This means that electron-hole pairs are created, and that the system is in an excited state. The excited electron can relax back due to electron-lattice scattering, in a characteristic time scale called the relaxation time $\tau$. Note that $\tau$ also describes the electron transport in a metal in the Drude model approximation [1]. From measurements of the temperature dependency of the conductivity of metals we know that $\tau$ decreases with temperature, so the Breathing surface model predicts a decreasing damping parameter with temperature.

Figure 2.4 shows the temperature dependence of the damping parameter for the 3d transition metals (note that the LL version of the damping parameter is plotted in this graph). We observe a steep decrease in damping from 0 K to approximately 100 K. This behavior is captured by the breathing Fermi surface model. However, for example for Co, we observe an increase in damping when the temperature is increased further. This is not captured by the breathing Fermi surface model. Because at room temperature, this is the dominant behavior, we need a more extended model to describe the intrinsic magnetic damping.

**Torque-correlation model**

One model that describes the whole temperature range is the torque-correlation model introduced by Kamersky [33]. This model gives the damping parameter as a function of the electron transition probability $|\Gamma_{nm}|$ between the electronic bands:

$$\lambda = \frac{g^2 \mu_B^2}{\hbar} \left[ \sum_{nm} \int \frac{d^3k}{(2\pi)^3} |\Gamma_{nm}^-|^2 \times W_{nm}(k) \right],$$

where $g$ is the Landé factor, $\mu_B$ the Bohr magneton, $\hbar$ Planck’s constant, and $W$ the spectral overlap of the electron spectral functions, which will be defined in a moment. We will not enter into details, but try to give a short explanation of the physical background behind this equation. $|\Gamma_{nm}^-|$ measures the transition of electrons between states in bands $n$ and $m$ induced by the spin orbit torque. There are two possibilities, $n=m$, which are called intraband transitions, and $n \neq m$ which is called an interband transition. Those transitions are weighted by the spectral overlap of the electron spectral functions $A_{nk}$:

$$W(k) = \frac{1}{\pi} \int d\varpi_1 \eta(\varpi_1) A_{nk}(\varpi_1) A_{mk}(\varpi_1).$$

As explained by Kambersky in [28], the overlap between the spectral functions is proportional to $\tau$ for the interband transitions, and proportional to $1/\tau$ for the intraband transitions. Gilmore used an effective field analogy to show that the intraband term is equal to the breathing Fermi surface model discussed earlier [18]. From the dependency of the interband term on $\tau$, we see that the interband contribution decreases...
2.2. MAGNETIZATION DYNAMICS

with increasing temperature and the intraband contributions increases with increasing temperature. So the intraband contribution is responsible for the observed increasing damping for temperatures above approximately 100 K in figure 2.4.

2.2.2.2 Enhanced damping by spin-pumping

When we take a look at some experimental values for the Gilbert damping parameter $\alpha$, we find for bulk Co a value of approximately 0.005 [34]. However in thin films like for example Pt-Co, the measured Gilbert damping parameter can be much larger, up to 0.22 (see for example chapter 5 of this thesis). Extrinsic mechanisms are believed to cause this increase in damping in thin films. Tserkovnyak et al [29] proposed a process called spin pumping to explain the enhanced damping in thin magnetic films in contact with another material. We will use a simplified picture to explain how spin pumping can lead to enhanced magnetic damping. We use figure 2.5 to explain the process. When the magnetization points along the effective magnetic field, the system is in equilibrium. On the left side of figure 2.5 we see a simplified density of states (DOS) of a 3d transition metal. Only the DOS for the 3d electrons, responsible for ferromagnetism, is represented. The unequal number of spin-up and spin-down electrons leads to a net magnetic moment in ferromagnets such as Co, Ni and Fe. Because the spin magnetic moment points anti-parallel to the spin, the magnetization will point in the direction of the minority spins. When the magnetization is out-of-equilibrium, the spins do not point along the axis of the field but in a direction with a certain angle to the effective field.

Every individual spin can be projected on a quantization axis which is not parallel to the direction of the spin, giving an majority and minority component along that quantization axis. When we align this quantization axis with the effective field, we will see that there are less majority electrons and more minority electrons in comparison with the equilibrium situation. This means that the system is in an excited state. The electrons can relax in different ways: one way is to scatter with the lattice, transferring their orbital momentum to the lattice system. This is equivalent to the process describing the intrinsic damping in the previous section. However, there is another possibility: the 3d electrons can also interact with conducting s electrons. In this process the d electron can transfer angular momentum to the s electron, and the d-electron system relaxes back to equilibrium. The conduction electrons can diffuse through the material and are also allowed to cross the boundary between the ferromagnetic layer and an adjacent non-magnetic layer. This way, angular momentum can be transported to the adjacent layer leading to a loss of orbital momentum in the magnetic system and therefore an enhanced damping.

In this adjacent layer, the spins can accumulate at the interface. In this case, there can be a spin current back into the ferromagnet, opposing the damping. However, when the spins decay or leave the interface fast enough, the non-magnetic layer can
be regarded as an ideal spin-sink. It is clear that spin-pumping is only effective if the area of the interface between magnetic and non-magnetic layer is large in comparison to the volume of the ferromagnetic layer. It turns out that the spin-pumping effect has a considerable effect for magnetic layers thinner than approximately 10 nm.

2.3 Current influenced magnetization dynamics

In this section we will introduce current influenced magnetization dynamics. We start with an introduction of spin transfer torque. The spin of an electron injected into a magnetic layer can exert a torque on the magnetization. Subsequently we will take a look at the two mechanisms that were proposed in literature to explain current assisted switching measurements done on Pt/Co/AlOx systems as we already discussed in the introduction chapter: the spin Hall effect (SHE) and Rashba spin-orbit coupling. We believe that those effects could possibly explain the de-pinning measurements done in this group on Pt/Co/Pt systems.
2.3 CURRENT INFLUENCED MAGNETIZATION DYNAMICS

2.3.1 Spin transfer torque

When an electron is injected into a ferromagnetic system, its spin can exert a torque on the magnetization. This torque is called spin transfer torque (STT) and is explained in this section. First, we consider a spin current injected into a ferromagnetic system. Secondly, we take a look at a current running through a ferromagnetic material with a gradient in the magnetization, for example a domain wall between two magnetic domains.

The total magnetic moment in a ferromagnet is linked to the total angular momentum \( \vec{L} \), which is the sum of the orbital and spin angular momentum of the electrons (if we ignore the spin of the nuclei). We define the angular momentum density in our magnetic system as \( \vec{S}_2 = \vec{L}/V \) (units: \( J_s/\text{m}^3 \)), with \( V \) the volume of the system. The magnetization is defined as the magnetic moment density, \( \vec{M} = \vec{\mu}/V \), with \( \vec{\mu} \) the total magnetic moment in the system. The magnetization can be written as a function of the total angular momentum density:

\[
\vec{M} = -\frac{g \mu_B}{\hbar} \vec{S}_2, \tag{2.22}
\]

with \( g \) the Landé factor, \( \mu_B \) the Bohr magneton and \( \hbar \) Planck’s constant.

We now consider a ferromagnetic system with uniform magnetization, represented by the vector \( \vec{M} \). Subsequently, we inject a spin current. A spin current is a current that carries a net angular momentum in a certain direction defined by the polarization vector \( \vec{P} \):

\[
P = \begin{pmatrix} P_x \\ P_y \\ P_z \end{pmatrix}. \tag{2.23}
\]

We will assume that the spin current is fully polarized, which means that the length of the polarization vector is 1. This spin current injects an amount of angular momentum into the system. When we divide the total injected angular momentum by the volume \( V \) of the system we get the injected angular momentum density \( \vec{S}_1 \). We will now take a look how those injected spins can exert a torque on the angular momentum of the ferromagnet, and thus via equation 2.22 on its magnetization.

We can project the injected spins, who are along their polarization vector, in a majority and minority component by projecting the polarization vector along any quantization axis we choose. This means that we write the spin state as a linear combination of majority (\( |+> \)) and minority (\( |-> \)) spin states with respect to this quantization axis:

\[
\chi = u_1 |+> + u_2 |->, \tag{2.24}
\]

in which \( u_1 \) and \( u_2 \) are complex numbers. Because of the normalization condition for quantum mechanical states, we have 3 free parameters to construct \( u_1 \) and \( u_2 \) which
means that equation 2.24 still contains the full 3 dimensional information about the direction of the spin.

When we take the direction of the localized angular momentum $\vec{S}_2$ in the ferromagnetic system as the quantization axis, the majority component will lie parallel to $\vec{S}_2$, and the minority component will lie anti-parallel to $\vec{S}_2$. In chapter 1 we already discussed that the minority and majority spins have different scattering probabilities. This means that the minority spins will be lost which corresponds to a rotation of $\vec{S}_1$ into the direction of $\vec{S}_2$. In figure 2.6 we see the direction of $\vec{S}_1/\vec{t}$. The total torque exerted on the injected spins is defined as:

$$ T_{STT} = V \frac{d\vec{S}_1}{dt}.$$  

(2.25)

From figure 2.6 we can construct the direction of this torque: $(\vec{S}_1 \times \vec{S}_2) \times \vec{S}_2$. When we assume the magnetic system to be a closed system after the injection of the spins, conservation of angular momentum dictates:

$$ \frac{d\vec{S}_1}{dt} = -\frac{d\vec{S}_2}{dt}. $$  

(2.26)

This means that the torque exerted on $\vec{S}_1$ will be in the direction of $\vec{S}_2 \times (\vec{S}_1 \times \vec{S}_2)$. The size of this torque scales with the injected angular momentum. Each injected spin carries an angular momentum of $\hbar/2$, and the number of injected electrons is given by $I/e$, in which $I$ is the injected current and $e$ is the charge of the electron. The exact size of the torque is however complicated to calculate. Slonczewski gives an expression for the size of this STT [7]:

$$ T_{STT} = g(\Theta) \frac{hI}{2eS_1S_2} \vec{S}_2 \times (\vec{S}_1 \times \vec{S}_2), $$  

(2.27)

in which $g = (\Theta)$ is a dimensionless scaling parameter which is a function of the angle $\Theta$ between $\vec{S}_1$ and $\vec{S}_2$. In systems comparable with ours, this factor $g$ is normally taken to be 1 [47].

When we want to add this STT to the LLG equation, we need to rewrite this expression into an equation which describes the evolution of the magnetization instead of the total angular momentum. By definition, we know that:

$$ \frac{d\vec{m}}{dt} = \gamma T_{STT}, $$  

with $\gamma = -\frac{e\mu_m}{2m_e}$ the gyromagnetic ratio. Further we define the spin direction as: $\vec{\sigma} = \vec{S}_1/|S_1|$. The direction of the magnetization can be written as $\vec{m} = -\vec{S}_2/|S_2|$. Combining those definitions and equations 2.22, 2.27 and 2.28, we get:
2.3. CURRENT INFLUENCED MAGNETIZATION DYNAMICS

Figure 2.6 – Direction of the torques exerted on each other by the magnetic spins $S_2$ and the injected spins $S_1$, as explained in the text. Remember that the magnetization direction is anti-parallel to the spin direction.

\[
\frac{d\vec{m}}{dt} = g(\Theta)\frac{hI\gamma}{2eVM}\vec{m} \times (\vec{\sigma} \times \vec{m}). \tag{2.29}
\]

When we define a magnetic system with length $L$, width $W$ and thickness $t_{FM}$ (see figure 2.9b), we can write the current density as $J = I/Wt_{FM}$ and the volume as $V = LWt_{FM}$ and we get:

\[
\frac{d\vec{m}}{dt} = g(\Theta)\frac{hJ\gamma}{2eVt_{FM}M_{sat}}\vec{m} \times (\vec{\sigma} \times \vec{m}). \tag{2.30}
\]

This expression can be inserted into the LLG equation.

We will now take a look at the STT exerted by a current running through a domain wall in a magnetic material. Due to spin depended scattering, a current that is send through a ferromagnetic material will get spin polarized. In a domain wall, there is a gradient in the magnetization and thus in the direction of the angular momentum. A spin flowing through this domain wall will feel this gradient and by the mechanism described above for the isolated magnetic system, there will be a STT exerted on the electron spin which tends to align the spin with the angular momentum. Because of the conservation of angular momentum, the electron spin will also exert a torque on the magnetization.

In figure 2.7 a schematic representation of a magnetic domain wall is depicted. The magnetization in the domain wall will change its direction $d\vec{m}$ over a distance $dx$. This means that the angular momentum $\vec{S}_2$ will change its direction by $-\vec{d}\vec{m}$ over this distance. We assume that the spins move adiabatically through the domain wall, and
Figure 2.7 – Schematic representation of adiabatic STT in a domain wall in a ferromagnetic wire. \( S_1 \) represents the angular momentum density of the conduction electrons, \( S_2 \) represent the angular momentum density of the localized \( d \) electrons. The conduction electrons will follow the angular momentum of the localized electrons, and their direction change is equal to \( -d\vec{m} \) over a distance \( dx \) as explained in the main text.

thus that the spin is always aligned with \( \vec{S}_2 \). We can express the change of angular momentum of the conduction electrons over the distance \( dx \) as the total amount of angular momentum flowing through the domain wall with length \( dx \) per unit of time multiplied with \( d\vec{m} \):

\[
\frac{d\vec{S}_1}{dt} = -P \frac{\hbar}{2} \frac{nv}{dx} d\vec{m}, \tag{2.31}
\]

in which \( v \) is the velocity of the electrons, \( P \) the polarization of the current, \( n \) the volume density of conduction electrons and \( \hbar/2 \) the angular momentum per electron.

When we use the definition of the electron current, \( J_e = n_e v e \), we will get:

\[
\frac{d\vec{S}_1}{dt} = -P \frac{\hbar}{2} \frac{J_e}{n_e dx} d\vec{m}. \tag{2.32}
\]

Using equation 2.26 and the relationship between magnetic moment and angular momentum, equation 2.22 we get:

\[
\frac{d\vec{M}}{dt} = -P g\mu_B J_e \frac{d\vec{m}}{dx}. \tag{2.33}
\]

This means that we can write the torque on the magnetization due to the spin polarized current (which we will call the adiabatic spin transfer torque \( T_{STT}^{adia} \) and using equation 2.28 for the definition of torque) as:

\[
T_{STT}^{adia} = -P g\mu_B J_e \frac{d\vec{m}}{2eM_{sat}\gamma dx}. \tag{2.34}
\]

We can generalize this equation to three dimensions:

\[
\frac{d\vec{m}}{dt} = -(\vec{u} \cdot \vec{\nabla})\vec{m}, \tag{2.35}
\]

with \( \vec{u} = P g\mu_B J_e \frac{\vec{J}_e}{2eM_{sat}} \) defined as the effective velocity.
2.3. CURRENT INFLUENCED MAGNETIZATION DYNAMICS

This expression can be added to the LLG equation to describe domain wall motion. However, experiments could not be fully explained by this equation. Therefore, a second STT term was added, called the non-adiabatic spin transfer torque. This torque is constructed in such a way that it is perpendicular to the adiabatic torque and the magnetization. The size of the torque is described phenomenologically with the non-adiabatically parameter $\beta$. This leads to the following expression for the non-adiabatic STT:

$$T_{\text{STT,non-adi}} = \frac{\beta}{\gamma} \vec{m} \times [(\vec{u} \cdot \vec{\nabla}) \vec{m}].$$ (2.36)

The exact origin of the non-adiabatic STT is still under debate [46]. It has to do with the fact that in samples with high gradients in the magnetization, the spins are not able to follow the magnetization perfectly like we assumed when we derived the adiabatic STT. This leads to a non-zero non-adiabatic STT. When we add both adiabatic and non-adiabatic STT terms to the LLG equation, we get:

$$\frac{d\vec{m}}{dt} = -\gamma G \mu_0 \vec{m} \times \vec{H}_{\text{eff}} - \alpha \vec{m} \times \frac{dm}{dt} - (\vec{u} \cdot \vec{\nabla}) \vec{m} + \beta \vec{m} \times [(\vec{u} \cdot \vec{\nabla}) \vec{m}].$$ (2.37)

The STT is able to drag a domain wall along with the current. The velocity of the domain wall as a function of the effective velocity is depicted in figure 2.8. When $\beta = 0$, we need a certain critical current to move the domain wall. When $\beta$ is not equal to zero, there is no critical current. The domain wall velocity as a function of the current is heavily dependent on the relative size of $\beta$ compared with $\alpha$. More information about domain wall motion can be found in [48]. Current induced domain wall motion allows for the development of new types of memory devices, like the racetrack memory we discussed in the introduction chapter.
2.3.2 Spin Hall effect

In this section we will introduce the spin Hall effect (SHE). When an electron is driven through a material with high spin orbit coupling, a transverse spin current is observed, see figure 2.9a. In the previous section, we discussed that when a spin current is injected into a ferromagnetic layer, it could exert a STT on the magnetization. This SHE induced STT torque could be a promising new technique to switch the magnetization in new memory devices. In this section we will give a phenomenological description of the SHE. We will also discuss why we need an in-plane magnetic field to be able to assist the magnetic switching using the SHE in perpendicular magnetized Pt/Co(B)/Pt nano-wires. We end this section by a discussion of the effect of the SHE induced STT on the magnetic damping.

Already in 1879, years before the discovery of the electron, Hall observed an effect that is now called the Ordinary Hall effect (OHE) [8]. He sent a current through a material and applied an external magnetic field perpendicular to it, and measured a voltage difference perpendicular to both the current and field direction [23]. The deflection of the electrons from the current direction is caused by the Lorentz force, which is felt by moving electrons in a magnetic field. In ferromagnets, besides this OHE, a second contribution comes into play called the Anomalous Hall effect (AHE), which is not proportional to the applied magnetic field but to the magnetization. Due to both intrinsic and extrinsic (spin dependent scattering) mechanisms, spin up and spin down electrons will deflect in opposite direction, perpendicular to the current direction. The spin dependent deflection leads to spin accumulation at the edges and, because of the imbalance between spin-up and spin-down electrons in ferromagnets, also to a charge accumulation. This leads to a measurable voltage difference transverse to the current direction. A Hall resistivity can be defined as the transverse electric field per unit longitudinal current density:

\[
\rho_H = R_0 H + R_s M,
\]

in which \( R_0 \) is the ordinary Hall coefficient and \( R_s \) is the anomalous Hall coefficient, which is experimentally found to be much larger than the ordinary Hall coefficient.

In a nonmagnetic metal, there is no imbalance between spin up and spin down electrons, but due to similar mechanisms that caused the AHE, electrons with opposed spin will still be deflected in opposite direction. Because the amount of spin up and spin down electrons is equal, there will be no charge separation, so no voltage difference over the system will appear. However, there will still be a spin imbalance. This effect is called the spin Hall effect (SHE). Although we do not go into further details about the origin of the SHE, we have to mention one important property: the spin dependent deflection originates from spin-orbit coupling. Therefore, the SHE is large in systems with a strong spin-orbit coupling. For more information about the origin of the SHE we refer to [17].
2.3. CURRENT INFLUENCED MAGNETIZATION DYNAMICS

![Diagram](image)

**Figure 2.9** – (a) Schematic representation of the spin Hall effect. Spin-dependent deflection of carriers produces a spin separation. For clarity, only spin separation in the $y$ direction is shown. The same effect takes place in the $z$ direction. The length, width and thickness of the system are labelled respectively $L$, $W$ and $t$. (b) Definition of a non-magnetic/magnetic/non-magnetic system.

When we consider an isolated non magnetic system with high spin-orbit coupling, the spin separation leads to a spin accumulation at the edges of the system. When we add an adjacent magnetic layer, the spins are able to diffuse into the magnetic layer. In this way we are able to inject a spin current into the ferromagnetic layer which is able to exert a STT on the magnetization like we discussed in the previous section. The spin current density entering the magnetic layer is phenomenological linked to the electron current density ($J_e$) through the (non-magnetic) material by a material constant called the Hall angle ($\Theta$) [15]:

$$\Theta_{SHE} = \frac{J_s}{J_e}. \quad (2.39)$$

As we mentioned before, the SHE is strong in systems with strong spin orbit coupling, so the Hall angle is material dependent. The Hall angle also depends on the non-magnetic layer thickness. The reason is depicted in figure 2.10b in which a Pt layer is show with a CoB layer on top. The spins drifting in the positive $z$ direction due to the SHE will get absorbed by the CoB layer, but the spins drifting down will accumulate at the edge of the sample. Those spins can diffuse back in the direction of the CoB layer and thereby reducing the effective spin current. Those diffusive back-flow of spins is limited by spin flip scattering. The typical distance a spin travels before its spin is flipped is described by the spin flip length $\lambda$, which depends on the material. If the Pt layer is made thicker, less spins diffuse back without losing their spin orientation,
and therefore the Hall angle is increased. The Hall angle as a function of the layer thickness for Pt is given in figure 2.10a.

Combining the definitions of the spin Hall angle with the Slonczewski torque from the previous section, we find a STT due to the SHE equal to:

\[
\frac{d\vec{m}}{dt} = g(\Theta) \frac{hJ_e \theta \gamma}{2et_{FM}M_{sat}} \vec{m} \times (\vec{\sigma} \times \vec{m}).
\] (2.40)

From figure 2.10 we know that the Hall angle for platinum is of the order of magnitude of 0.05. Despite this small number, the SHE allows us to inject large spin-currents into the ferromagnetic layer, due to the ratio between the areas through which the electron and spin currents are flowing. This becomes clear when we calculate the total spin current as a function of the total electron current in a system with length \(L\), platinum layer thickness \(t\) and width \(W\) (see figure 2.9a):

\[
I_s = \frac{L}{t_1} I_e \Theta.
\] (2.41)

Some typical numbers used in this research are: \(L = 1\) mm and \(t_1 = 5\) nm and \(W = 50\mu m\) and \(\Theta = 0.05\). For these values we find:

\[
I_s = 10^4 \cdot I_e.
\] (2.42)

This means that we can use the SHE to inject large spin currents into our magnetic layer. Those large spin currents could be a great tool to switch the magnetization of the magnetic layer. However, to assist magnetic switching, the SHE induced STT should break the symmetry of the up and down magnetization direction. A symmetry
2.3 CURRENT INFLUENCED MAGNETIZATION DYNAMICS

Figure 2.11 – The precessional motion of the magnetization around respectively the x, y and z axis is depicted. Also the STT induced by the SHE in Pt/Co(B) nano-wire with an applied current in the x direction is depicted. The direction of the injected spins is indicated by $\sigma$.

analysis given by Haazen [17] shows that the SHE induced STT on its own is not able to destabilize one of the two magnetization directions. We need an in-plane magnetic field in the x direction to break the symmetry and introduce a preferred direction of the magnetization due to the SHE induced STT.

We will now take a look at the influence of the SHE on magnetic damping. We explain why the direction of the in-plane field is important for the size of the effect of the SHE induced STT on the magnetic damping. In figure 2.11, the precessional motion of the magnetization in a magnetic nano-wire around respectively the x, y and z axis is depicted. When we apply a current in the x direction, we inject spins which are aligned along the y direction. The effect of those injected spins is that it tends to pull the magnetization in along the y axis, as we explained earlier. The direction of this torque is given by equation 2.40. When the magnetization processes around the y axis, the SHE induced STT is pointing to the center of the precessional motion during the whole cycle, or in the opposite direction depending on the current direction. The injected spins want to align the magnetization with the y axis, and are therefore assisting the damping. This also becomes clear when the direction of the torque is evaluated using 2.40. From this evaluation it becomes clear that this torque is always parallel or anti-parallel to the Gilbert damping term, and thereby the effect on the damping is maximized. When the magnetization is precessing around either the z or x axis, the STT will not point toward the center of the precession, but always roughly in the same direction, see figure 2.11. Therefore, during half of the cycle, the damping is
assisted, and during the other half the damping is opposed. Therefore, the net effect is minimized for those configurations. When the magnetization is precessing around an effective field that is in an arbitrary direction, the effect on the damping will have some intermediate value.

2.3.3 Rashba spin-orbit coupling

In this section we will discuss the second spin-orbit induced torque we introduced in the introduction chapter, the Rashba spin orbit-coupling which is induced when a current is flowing at the interface between a material with high spin-orbit coupling and a second layer of a different material. The effect was first described by Rashba for semiconductors [36], but is also present in metallic systems with high spin-orbit coupling. Electrons at the interface between two materials experience an electric field \( E_{\text{s,i.a}} \), in which \( \vec{e}_z \) is a unit vector in the z direction and “s.i.a” stands for “spin inversion asymmetry”. This electric field originates from the different ionic potentials of the atoms of the two materials. In a classical picture it is easy to see how this lead to a magnetic field. An electron moving through an electric field feels a magnetic field. A Lorentz transformation will give us the direction of this so called Rashba field, which will be perpendicular to both the electric field and the current direction.

Quantum mechanically the Rashba field originates from spin-orbit coupling introduced in section 2.1.1. When the s.i.a electric field is inserted in equation 2.13 and assuming no applied magnetic fields (\( A=0 \)) we get:

\[
H = \frac{1}{2m_e}(\hat{p})^2 + V - \frac{\hbar^2 E_{\text{sia}}}{4m_e^2 c^2} \vec{\sigma} \cdot (\vec{k} \times \vec{e}_z). \tag{2.43}
\]

This last term looks just like the energy of a magnetic moment in a magnetic field, given by the Zeeman energy, \( H_{\text{zeeman}} = \mu_0 (\vec{M} \cdot \vec{H}) \). This means that the spin-orbit term of the Hamiltonian leads to an effective field effect, which we will define as the Rashba field \( H_R \):

\[
\vec{H}_R \sim \vec{\sigma} \cdot (\vec{k} \times \vec{e}_z). \tag{2.44}
\]

This effective field is in principle only felt by the conduction electrons, but couples via s-d interaction to the local magnetization. The size of this effective magnetic field is calculated by Manchon and Zhang [26]and given by:

\[
\vec{H}_R = -2 \frac{\alpha_R m}{\hbar e M_s} P(\vec{J}_C \times \vec{e}_z), \tag{2.45}
\]

in which \( J_C \) is the charge current density, \( P \) the polarization of the electron current, \( M_s \) the saturation magnetization. The strength of the spin-orbit coupling is phenomenological taken into account by the parameter \( \alpha_R \), which is proportional to \( E_{\text{sia}} \).
2.4. SIMULATIONS OF CURRENT INDUCED MAGNETIZATION DYNAMICS

In our systems, the Rashba field would lie in the y direction, perpendicular to both the current direction and the out-of-plane axis of the sample, see figure 2.12.

This Rashba field exerts a torque on the magnetization given by $\vec{T}_R = \vec{m} \times \vec{H}_R$. However, this term could not explain the switching and depinning measurements because it is not able to break the symmetry of the two magnetization directions. However, in literature a second non-adiabatic contribution from the Rashba spin-orbit has been proposed: [16]:

$$\vec{T}_{R,na} = \alpha_{R,na} \vec{m} \times (\vec{H}_R \times \vec{m}),$$

in which $\alpha_{R,na}$ is the non-adiabatic Rashba parameter. This torque is in the same direction as the STT from the SHE and in combination with an applied field in the x direction has the right symmetry to explain the switching experiments.

In systems with two interfaces the Rashba field from both interfaces is opposite, because $E_{sia}$ is in the opposite direction at the two interfaces. When the interfaces are identical, like in Pt/Co(B)/Pt, we expect those two contributions to the Rashba field to cancel each other. This was confirmed by Haazen for Pt/C/Pt [17]. Due to differences in current densities in the top and bottom platinum layer and differences in growth quality, there can still be a small Rashba field present.

2.4 Simulations of current induced magnetization dynamics

In the previous section we introduced current induced spin-orbit torques. The main focus in this research will be on Pt/Co(B)/Pt systems, for which we can neglect the Rashba contribution. In this section we discuss macro-spin simulations to investigate the effect of a SHE induced STT on the magnetization dynamics. In the first section, we discuss current assisted switching simulations which will demonstrate that in the proper configuration the STT induced by the SHE can oppose or assist the magnetic switching depending on the current and applied in-plane field direction. In the second section, simulations will demonstrate that the SHE induced STT is able to influence the magnetic damping. This means we can study the SHE by measuring the effect on the damping parameter.

During all the simulations, the saturation magnetization of Co is used instead that of CoB, which is used in the sample fabrication. The reason is that exact saturation
magnetization for our CoB used is not known but is close to that of Co.

In the simulations, we used values for the Hall angle reported for Pt, see figure 2.10 [15]. For the current distribution in the nano-wire we assume that all the current will flow through the two Pt layers. Lavrijsen [27] showed that the current density in the two Pt layers in asymmetric Pt/Co/Pt nano-wires is not equal. There will be a ratio between the current densities in the top and bottom layer:

\[
\frac{J_{\text{top}}}{J_{\text{bottom}}} = x. \tag{2.47}
\]

Lavrijsen found a value of \(x=0.93\) for Pt(4)/Co/Pt(2) nano-wires. We decided to use the same value for our Pt(5)/Co/Pt(2) system. This is a rough estimate but because the difference in current densities is too small to be the dominating effect, and is negligible in comparison with the difference in Hall angle for the two layers, we believe that this assumption is allowed.

### 2.4.1 SHE induced magnetic switching simulations

In this section simulations of current assisted field induced magnetic switching in a perpendicular magnetized 20 \(\mu\)m wide Pt(5)/Co(0.7)/Pt(2) nano-wires are discussed. We used a saturation magnetization of \(1.422 \times 10^6\) A/m and an anisotropy field of \(H_{\text{ani}} = -1.1 \cdot M_z A/m\). The simulations are performed in the following way. A field is applied under an angle of 75° with the z axis, see figure 2.13. The simulations are performed in two different configurations: with the field in the x-z plane and the y-z plane. Those two field configurations are labeled respectively \(H_{x-z}\) field and \(H_{y-z}\) field, to indicate the plane in which the field is applied. This is important, because the direction of the in-plane component of the field is important for the effect of the SHE induced torque.

The simulation starts at a large negative field, typically \(-400\) mT. The equilibrium position of the magnetization is calculated. Then the field is decreased in small steps of typically 0.5 mT, and every time the new equilibrium position is calculated. This is repeated until the field is 0, and then the field is increased with the same steps in the opposite direction. At some point the magnetization will switch. The field at which this happens is defined as the switching field \(H_{\text{switch}}\). Without applied current, the switching field obtained from the simulations is 110 mT. This is larger then is observed switching field in measurements, because the simulations do not take thermal nucleation of domains into account, which can grow due to domain wall motion and thereby flip the magnetization at lower fields.

Subsequently the LLG equation is extended with the STT originating from the SHE, see equation 2.40, and the simulations are repeated for different values of the current. A change in switching field is observed which is defined as \(\Delta H_{\text{switch}}\). The results for an applied \(H_{x-y}\) field are showed in figure 2.14. Simulations are performed for different values of the Hall angle, because the actual size of the Hall angle in our
2.4. SIMULATIONS OF CURRENT INDUCED MAGNETIZATION DYNAMICS

![Diagram of a nano-wire with H_yz and H_xz fields](image)

**Figure 2.13** – Definition of $H_{yz}$ and $H_{xz}$ for a wire lying along the $x$ direction.

**Figure 2.14** – Simulations of the change in switching field due to the SHE, in the $H_{xz}$ field configuration. The simulated system is a Pt(5)/CoB(0.7)/Pt(2) nano-wire. Simulations are performed with different values of the Hall angle. The percentage indicates the fraction of the Hall angle from literature [15] that is used.

From those simulations we can conclude that the switching is indeed assisted (or opposed) by the SHE. The effect is inverted when the current direction is inverted. Simulations with an applied $H_{yz}$ field are also performed. In this case, there was no effect visible on the switching field. So a field in the $x$ direction is necessary to assist or oppose the switching field with the SHE, as is expected from symmetry considerations mentioned in section 2.3.2.
2.4.2 SHE influenced magnetic damping simulations

In this section, macro-spin simulations used to investigate the influence of a SHE induced STT on the magnetic damping are discussed. The Hall angle estimated by comparing current assisted switching experiments and simulations, as explained in chapter 4, is used. An intrinsic Gilbert damping parameter of 0.2, an external applied magnetic field of 500 mT and an angle of 80 degrees with the z axis are used in this section. The magnetization dynamics for the first 900 ps after the system is brought out of equilibrium is simulated. As discussed before, the magnetization starts a precessional motion that will damp out to its equilibrium position. A typical trajectory of the z component of the magnetization is given in figure 2.15. The measured or simulated trajectory of the z component of the magnetization is fitted with a phenomenological function [40]:

\[
\frac{\Delta M_z}{M_z} = A_1 \cdot +A_2 \cdot e^{-\frac{t}{t_d}}\sin(\omega t + \phi_0),
\]  

(2.48)

in which \(A_1\) and \(A_2\) are scaling parameters, \(t_d\) is the electron-phonon relaxation time, \(\omega\) is the precession frequency and \(\phi_0\) is the the phase of the oscillation. When we want to fit the complete trajectory, including the re-magnetization process after demagnetization, we can use an extended version of this equation, as described by Jozsa [40]. Kuiper [8] solved the LLG equation for a perpendicular magnetized sample with the help of LLG simulations, and found a very useful relation between the Gilbert damping parameter and \(t_d\) and \(\omega\):

\[
\alpha = \frac{1}{t_d\omega}.
\]  

(2.49)

When we obtain \(t_d\) and \(\omega\) from the trajectory of \(M_z\) using the fit function, this equation gives us the effective Gilbert damping parameter \(\alpha\).

Results from two types of simulations will be discussed in this section. Figure 2.16 shows results for simulations on a Pt(5)/Co(0.7)/Pt(2) wire for both an applied \(H_{xz}\) and \(H_{yz}\) field. For the \(H_{xz}\) field configuration an effect on \(\alpha\) is observed, but only for very high current densities. On top of that, the sign of the effect is not affected by the inversion of the current direction. For the \(H_{yz}\) field configuration, a much larger effect at much lower currents is observed. On top of that, the effects changes sign when the current direction is inverted. For these reasons we will focus on the \(H_{yz}\) field configuration for the rest of the simulations and for the measurements.

The second series of simulations shows the effect of the Pt layer thickness. We showed earlier that the Hall angle depends on the Pt thickness. From the simulations we determine the change of \(\alpha\) per unit of applied current. The results are shown in figure 2.17. The results are plotted in two ways: the change in \(\alpha\) as a function of the current density, and the change in \(\alpha\) as a function of the total current. The first is plotted in figure 2.17a. The observed effect increases for increasing layer thickness.
2.4. SIMULATIONS OF CURRENT INDUCED MAGNETIZATION DYNAMICS

Figure 2.15 – Plot of the z component of the magnetization simulated using the LLG equation.

Figure 2.16 – Simulations of the effective Gilbert damping parameter as a function of the current density. Simulations are performed in both the $H_{yz}$ and in the $H_{xz}$ configuration.
and it saturates for thick layers. The latter is plotted in figure 2.17b. In the first part of the plot the same behavior is observed: the effect increases with increasing layer thickness. But at some point the effect starts to decrease when the layer thickness is further increased. This is due to the fact that increasing the layer thickness further does not increase the Hall angle much more and higher current are necessary to obtain the same current density.

The reason the second graph is plotted, is Joule heating. Joule heating limits the amount of current that can be send though the nano-wires. Heat diffusion through the substrate scales with the area of the wire, thus Joule heating is dominated by the total current and not by the current density. Therefore, the total current that can be send through the wire is limited, not the current density. Figure 2.17b shows that the effect is optimized for a given current for a bottom platinum layer thickness of 5 nm.
Chapter 3

Methods

In this chapter the sample fabrication and measurement techniques used in this research to study current assisted magnetization dynamics are discussed. In the first section of this chapter we will take a look at the sample fabrication of thin films and magnetic nano-wires. For the deposition, magnetron sputtering is used, which gives the opportunity to deposit layers with sub-nanometer precision. Electron beam lithography (EBL) is used to pattern designs on a Si substrate. A last technique used during sample fabrication is plasma oxidation in combination with annealing to create AlOx layers. In the last part of the first section, details about the samples fabricated for this research are discussed.

The second section of this chapter describes the two measurement techniques used during this research. We are interested in studying ultrafast magnetization dynamics on the ps timescale. A very suitable method to study magnetization dynamics on the ps timescale is the time-resolved MOKE set-up. Also an imaging technique able to provide a spatial map with a 100 nm resolution of the magnetic state of a sample, the Kerr microscope, is introduced.

3.1 Sample fabrication and design

Both simple thin films as well as nano-wires with micrometer scaled features are used in this research. The production of those samples involves a couple of steps which will be discussed in this section. Also the design and materials used for the systems studied in this thesis are discussed.

3.1.1 Magnetron sputtering for thin film deposition

The samples are grown using a DC magnetron sputtering process, with the CARUSO ultra-high vacuum (UHV) system available in this group. The system consists of a vacuum vessel, in which the substrate is loaded. The base pressure inside the vessel is
around $3 \times 10^{-8}$ mbar. Inside the vessel there is a sample stage and six target materials placed above it. The sample stage is rotatable which allows the sample to be placed below the desired target.

When the sample stage carrying the sample is placed below the desired target, a small amount of Ar gas is added to the chamber, increasing the pressure to roughly $3 \times 10^{-2}$ mbar. Subsequently, a high voltage (100-1000 V) is applied between an anode and the target, which acts as a cathode. The Ar gas will get ionized and the ions are accelerated towards the target. When the ions hit the target they can knock single atoms out of the material, which will, due to diffusion, reach the substrate and be deposited on it. This process is schematically illustrated in figure 3.1a. The speed of this growing procedure is in the order of Ångströms per second, which means that deposited layer thickness can be controlled on the sub-nanometer scale.

### 3.1.2 Sample patterning

Using the magnetron sputtering process discussed above it is possible to control the deposited layer thickness which allows for the creation of simple thin-film samples. However, we want to be able to design more complex structures. We are interested in two categories of samples: Wedge films and patterned samples. The first type of samples are fabricated using the wedge mask that is present in the CARUSO facility. The patterned samples are fabricated using both a shadow mask and electron beam lithography (EBL), which allows us to basically fabricate any design we want with
3.1. SAMPLE FABRICATION AND DESIGN

features even below the micrometer scale. The use of masks and EBL are discussed in the following two subsections.

3.1.2.1 Masks

A couple of masks are available in the CARUSO system. Those masks are metal plates with cut-out designs which can be placed just above the substrate. We used two of those masks. The first is the wedge mask, which allows for the deposition of thin layers which vary in thickness in one lateral direction. The mask basically consists of a shutter which covers partly the substrate, and can be displaced to cover more or less of the substrate. It is placed 7.4 mm above the sample holder. When the mask is moved slowly during deposition, a gradient in the layer thickness is created as illustrated in figure 3.1b.

The second mask is used when we fabricate nano-wires. We want to apply a current through those nano-wires, and therefore the samples have to be placed into a chipcarrier, which can be connected to a power source. To be able to connect the micrometer scaled nano-wires to the chipcarrier, we need the nanowires to be in contact with large contacts on the sample which can be connected with the chipcarrier using a wire bonder. Those large contacts would be very time consuming to fabricate using EBL. Therefore a shadow mask is used containing a pattern of flower-like structures, each containing eight of those large contacts, as depicted in the left side of figure 3.4. Subsequently, EBL is used to pattern nano-wires in contact with those contacts as will be explained in the next section.

3.1.2.2 Electron Beam lithography

For our research we want to fabricate nano-wires with a length of about 1 mm and a width of 20-50 μm. Electron beam lithography (EBL) is an excellent technique capable of patterning samples with μm scaled features. The process is illustrated in figure 3.2. First, two layers of the photo-resist material Polymethylmethacrylaat (PMMA) are spin-coated on the substrate. Spin coating is done at 5000 rpm for 50 seconds, and is followed by a baking step at 150 °C for 120 seconds. During this baking step the solvent (anisole) is removed. The first layer consists of smaller polymers (495 repeat units) and is therefore softer then the second layer which consists of longer polymers (495 repeat units). Because the first layer is softer, an undercut is created after the development step, which leads to better edges in the samples.

Subsequently, the sample is loaded into the FEI Nova 600i Nanolab dual beam system which is used for the EBL step. The dual beam system is able to irradiate samples with either an electron beam or an ion beam. We use Raith software to control the electron beam and irradiate the desired path on our sample. Typically we use electrons with an energy of 30 keV and an dose of 300 μC/m² in this research. Because
CHAPTER 3. METHODS

Figure 3.2 – Schematic representation of the sample fabrication using Electron Beam Lithography (EBL).

3.1.3 Oxidation

Besides Pt/CoB/Pt samples, we are also interested in studying samples for which the upper Pt layer is replaced by an Al layer. The reason is discussed in chapter 4. In chapter 2 we discussed that perpendicular anisotropy in Pt/CoB/Pt systems is caused by the strong bonding between Pt and Co atoms at the interface. The bonding between Co and Al however is not that strong. To obtain perpendicular anisotropy, the Al has to be oxidized. In that case the Co atoms form a strong bonding with the O atoms in the AlOx layer which leads to perpendicular anisotropy. Although natural oxidation can oxidize Al up to approximately 1.5 nm, this process is slow (15 hours to oxidize 1nm [39]) and less controllable, and therefore we choose to use plasma oxidation in the oxidation chamber attached to CARUSO to oxidize the Al in a controlled way. The base pressure in the oxidation chamber is $1 \times 10^{-9}$ mbar. Oxidation takes place in a pure oxygen environment at a pressure of $1 \times 10^{-1}$ mbar. A current of 7.5 mA is send
through the oxygen in order to ionize it. The oxidation is followed by a baking step which is performed in an Ar environment at a temperature of 250 °C for 30 minutes. This baking step is necessary to create a nice AlOx structure by diffusing the oxygen homogeneously through the Al.

It is important to oxidize the Al long enough to oxidize the complete layer, but not too long because then the CoB starts to oxidize. Therefore the oxidation time should be optimized. However, it is easier to fix the oxidation time and tune the Al thickness. We fabricated a Pt/CoB/Al sample with an Al wedge, see figure 3.3a. This sample is oxidized using the above explained procedure for 320 seconds and subsequently $M_z$ as a function of Al thickness is measured. The results are given in figure 3.3 b. When we use an oxidation time of 320 seconds, the system is out-of-plane magnetized for a Al layer thickness of 0.95 nm.

### 3.1.4 Sample design

In this last subsection, the design and materials used for the samples in this thesis are discussed, starting with the choice of the magnetic layer. In the beginning of this research we were fabricating samples with magnetic layers made of Co, in order to stick as close as possible to previous work in this group. However, at some point we tried CoB because previous research showed that CoB has lower anisotropy [27] and therefore could have a lower damping parameter. When the damping is lower, more oscillations can be observed which makes the determination of the damping parameter more accurate. Although we did not observe a measurable difference in damping, the
magnetic signal increased significantly which leads to a better signal to noise ratio. For this reason we decided to use CoB from that moment on.

Also the choice of the substrate turned out to be very important when performing current induced measurements in the TR-MOKE set-up. The samples are fabricated on Si wafers. We started with wafers with a thin 10 nm oxide layer (only natural oxidation), because SiOx is a much better thermal insulator than pure Si. Because heating becomes an issue when high electrical currents are send through magnetic nano-wires in combination with a high power pump-beam, we decided to use this substrate. However, this substrate is not suitable for electric measurements due to a strange solar-cell effect. When nano-wires were fabricated on this substrate in a later stage of the research, a strange behavior was observed during current assisted damping measurements. We measured the current through the wire for a certain applied voltage. A dependency of the measured current on the presence of the pump-beam and and to a lesser extent the probe-beam was observed. For example, when a very small voltage of 0.005 mV is applied, a current of 0.0003 mA is measured. With both the pump-beam and probe beam focused on the wire, we measured a current of 0.0136 mA, a huge difference. This can not be explained by a change in resistance of the wire as a consequence of a temperature increase due to the pump beam, because the sign of the effect depends on the current direction. It seems that there might be a photovoltaic or a thermoelectric effect at work in our sample [43], but more research is needed to explain the physics behind this effect.

To exclude this photovoltaic or thermoelectric from happening we decided to use Si wafers with a 100 nm SiOx layer on top. This thick SiOx layer isolates the wires electrically from the Si. We verified that in those samples we do not observe the solar-cell effect, and that the heating of the sample is not a problem when not too high currents are used. As will be discussed in chapter 5, the signal to noise ratio turns out to be much better on this substrate.

Figure 3.4 – a) Schematic representation of a sample containing four ‘flowers’ which is placed in a chipcarrier. The contacts on the sample are connected to the chipcarrier using a wire bonding technique. b) An enlargement of one flower, containing two nanowires. The nano-wires are sputtered in contact with the flowers, in order to connect them to the chipcarrier. c) Enlargement of one nano-wire in contact with the flower contacts.
3.1. SAMPLE FABRICATION AND DESIGN

The design of our wires is given in figure 3.4. The shadow mask is used to fabricate the flower shaped contacts depicted in figure 3.4a, which are made of 30 nm Pt or Au. Each flower consists of eight contacts which point to the center of the flower. The size of a sample containing four flowers is such that it fits in the sample carrier we use as depicted in figure 3.4a. The flowers are connected with contacts on the chip-carrier with gold wires using a wire bonding technique. In figure 3.4b a magnification of the center of one of the flowers is depicted. In this area the magnetic nano-wires are fabricated in contact with the flower. In figure 3.4c, an enlargement of one of the nano-wires is depicted.

When thin wire is sputtered on a thick contact which has a sharp edge, we face the problem depicted in figure 3.5. When the wire is sputtered in contact with a contact that is not too thick, there is no problem, see the top image. However, when the contact is too thick, the nano-wire could break and this leads to a bad or no electrical contact. This problem could be solved by first sputtering the nano-wire next to the contact, and in a second step sputtering a thick layer of a conducting metal on top of both the contact and the nano-wire to connect them as depicted in figure 3.6a. However, the flower shaped contacts sputtered in CARUSO have soft edges, which makes it possible to sputter the nano-wires directly on the contacts, see figure 3.6b. This saves one EBL and one sputtering step. For this reason, we used this last technique for most of the samples fabricated in this research.

In this thesis both Pt/CoB/Pt and Pt/CoB/AlOx samples are fabricated. The oxidation process for the Pt/CoB/AlOx is the same as for thin films, as described in section 3.1.3. However, the oxidation time had to be re-calibrated, probably because the combination of oxidation time and Al thickness was determined using a sample that was fabricated using the wedge mask. It turns out that the calibration of the layer thickness as a function of time is not completely accurate when the mask is used. Therefore, in the nano-wires an Al layer thickness of 0.89 nm was used to obtain out-of-plane anisotropy.
Normal procedure

1 30nm Flower (pt/Au)

2 30nm Flower (pt/Au)

3 30nm Flower (pt/Au)

Optimized procedure

1 30nm Flower (pt/Au)

2 30nm Flower (pt/Au)

(a)

(b)

Figure 3.6 – a) The solution to the problem depicted in figure 3.5, is to first sputter the wire next to the contact, and then sputter a thick conducting material on top of both the flower and the wire to make a good contact. b) When the edge of the flower is smooth, it is possible to sputter the wire directly in contact with the flower and still obtain a good electrical contact.

3.2 Measurement methods

In this section two measurement set-ups used during this research are discussed. Both are optical set-ups using the so called magneto optical Kerr effect (MOKE). Therefore we start this section with an explanation of this effect in the first subsection. To study the magnetization dynamics the magnetization has to be brought out of equilibrium. This is done by demagnetize the system using a powerful laser pulse. This laser induced demagnetization is introduced in section 3.2.2. Subsequently the details of the Time-Resolved MOKE setup and the Kerr-Microscope are discussed.

3.2.1 Magneto optical Kerr effect

The magneto optical Kerr effect (MOKE) is closely related to another optical effect called Faraday rotation or Faraday effect [38]. The Faraday effect was discovered by Michael Faraday in 1845 and was the first indication that light and electromagnetism are related. When a polarized electromagnetic wave travels through a magnetic field, its polarization axis will tilt over an angle which is proportional to the effective field component in the direction of propagation of the electromagnetic wave.

When an electromagnetic wave reflects on a magnetic material, something similar to the Faraday effects happens. When light interacts with a medium, its electric field component $\vec{E}$ will induce a displacement field inside the material: $\vec{D} = \bar{\epsilon}\vec{E}$, where $\bar{\epsilon}$ is the dielectric constant, which is in general a second rank tensor. Its diagonal terms are independent of the magnetization, but its off-diagonal terms are related to the magnetization inside the medium [10]. For the specific case of light traveling in the
3.2. MEASUREMENT METHODS

z-direction and a magnetic sample with perpendicular magnetization, the dielectric constant can be written as:

\[ \bar{\epsilon} = \begin{bmatrix} \epsilon_{xx} & \epsilon_{xy} & 0 \\ -\epsilon_{xy} & \epsilon_{xx} & 0 \\ 0 & 0 & \epsilon_{xx} \end{bmatrix}. \] (3.1)

Those magnetization dependent off-diagonal terms mix the E components of the electromagnetic field, thereby varying its polarization. This dielectric constant can be inserted into the Fresnel equation. The Fresnel equation follows from the Maxwell equations and describes the behavior of light when it crosses an interface between two materials with different refractive indices $n$. This leads to the following relation:

\[
\begin{bmatrix}
\epsilon_{xx} - n^2 & \epsilon_{xy} & 0 \\
-\epsilon_{xy} & \epsilon_{xx} - n^2 & 0 \\
0 & 0 & \epsilon_{xx} - n^2
\end{bmatrix}
\begin{bmatrix}
E_x \\
E_y \\
E_z
\end{bmatrix} = 0.
\] (3.2)

The polarization can be expressed using the so-called Jones formalism. The Jones formalism describes the polarization of light by formulating a Jones vector which describes the direction of its electric field component. The length of those vectors is normalized. For light traveling in the z direction, the polarization is described by a vector $(E_x \ E_y \ 0)^T$. For example, light polarized linearly in the x direction is represented by $(1 \ 0 \ 0)^T$, and right and left circular polarized light by $\frac{1}{\sqrt{2}}(1 \ i \ 0)^T$ and $\frac{1}{\sqrt{2}}(1 \ -i \ 0)^T$. The interaction between light and optical components can be described with Jones matrices, which will be used in subsection 3.2.3.

Equation 3.2 describes the reflection and transmission of electromagnetic waves that come across an interface of two different materials. We find two eigenvectors for the electric field inside the material, with two corresponding eigenvalues, which represent the refractive index:

$E_0 = (1 \ i \ 0)^T$, with eigenvalue $n^2 = \epsilon_{xx} + i\epsilon_{xy}$ and $E_0 = (1 \ -i \ 0)^T$, with eigenvalue $n^2 = \epsilon_{xx} - i\epsilon_{xy}$. Those eigenvectors represent right and left circular polarized light, which both have a different refractive index given by the corresponding eigenvalue.

Linear polarized light can be written as the sum of right and left polarized light (easy to see, for example, for x polarized light: $(1 \ i \ 0)^T + (1 \ -i \ 0)^T = (1 \ 0 \ 0)^T$). Those two components of linear polarized light will have different refractive indexes. This means that the linear polarized light will undergo a complex rotation:

\[ \Theta = \theta + i\varepsilon = \frac{\epsilon_{xy}}{\sqrt{\epsilon_{xx}(\epsilon_{xy} - 1)}}, \] (3.3)

where $\theta$ is defined as the Kerr rotation and $\varepsilon$ as the Kerr ellipticity. Because the off-diagonal components of the dielectric constants depend on the magnetization, measuring the Kerr rotation or the Kerr ellipticity of the reflected light gives information about the magnetization in the material.
CHAPTER 3. METHODS

Figure 3.7 – The three types of MOKE as explained in the text.

Until now we analyzed the situation for perpendicularly incident light on perpendicular magnetized samples. This is usually referred to as the polar MOKE. There are two more configurations, called longitudinal and transverse MOKE. In longitudinal MOKE, the light is incident under an angle with respect to the sample plane, and the magnetization is in-plane in the longitudinal direction. In transverse MOKE the light is again incident under an angle, but now the magnetization is in the transverse direction. Figure 3.7 shows the three types of MOKE.

When the magnetization is in an arbitrary direction, the Kerr rotation will be more or less sensible to the magnetization component along the polar, longitudinal and transverse direction, depending on the angle of incidence of the light.

This means that when perpendicular incident light is used on a sample with arbitrary magnetization direction, the Kerr rotation is most sensible to the out-of-plane component of the magnetization.

3.2.2 Ultrafast demagnetization

To study the precessional motion of the magnetization we have to bring the magnetization out of equilibrium. In chapter 2 we discussed that the effective field along which the magnetization tends to align itself consists of the applied magnetic field and an anisotropy field which depends on the magnetization. So to bring the system out of equilibrium, an external field pulse could be used. Another opportunity is to change the anisotropy field by demagnetizing the sample. This can be done by hitting the sample with an intense laser pulse. This demagnetization happens on a very short time scale and is therefore called ultrafast demagnetization. Its exact origin has challenged scientists for the last few decades.

In the 90’s, experiments gave an estimation of the demagnetization time in Ga of 100 +/- 80 ps [12]. Theoretical estimations of spin relaxation confirmed those numbers shortly after. Later, Beaurepaire et al [13] performed demagnetization measurements on Ni, and found demagnetization times well below 1 ps, so on a completely different timescale as found for Ga. Beaurepaire introduced a phenomenological three temperature model (3TM) to explain their measurements. In the 3TM model, the system is decomposed in three subsystems: the electron, phonon and spin system. A temperature and a heat capacity are assigned to every sub-system, and interactions between
3.2. MEASUREMENT METHODS

them are described by coupling constants. When the laser hits the sample, initially, only the electron temperature will rise on a timescale of 50-100 ps. Due to interactions, also the spin system will get thermalized, leading to a loss in magnetization. The spin system will eventually cool down due to interactions with the phonon system, leading to re-magnetization. This model is however not able to explain the different time-scales found in Ni and Co on one side and gadolinium on the other side. Koopmans et al. [11] implemented a microscopic description of spin-flip scattering into the 3TM, and they were able to describe the demagnetization in various materials with one single model, called the microscopic 3 temperature model (M3TM).

Ultrafast demagnetization can be used to bring the magnetization out of equilibrium in the following way. When no external magnetic field is applied, the effective field will lie along the easy axis and the magnetization will point in that direction. When the sample is demagnetized, the magnetization will be decreased but its directions will not change. However, when we apply a magnetic field, the magnetization lies along the effective magnetic field which is now the sum of the applied magnetic field and the anisotropy field, which depends on the size of the magnetization, see equation 2.11 in combination with equation 2.4. If we now demagnetize the sample, the direction of the effective magnetic field is changed, and this will bring the system out of equilibrium. This is used to measure the magnetization dynamics in the TR-MOKE set-up which is discussed in the next subsection.

3.2.3 TR-MOKE Set-up

In this subsection the time resolved (TR) MOKE set-up is discussed. The technique is based on the MOKE and uses a pulsed laser in combination with a stroboscopic measurements scheme. The pulsed laser is split into two beams, one to excite the system by demagnetize the sample (the pump-beam) and one to probe the magnetization using the MOKE (the probe-beam). By changing the delay between the pump and the probe-beam we are able to reconstruct the trajectory of the magnetization as a function of time on the ps timescale. In this section the set-up is explained in more detail.

A schematic representation of the TR-MOKE set-up is given in figure 3.8. A Tsunami mode-locked Ti:Sapphire laser from Spectra Physics is used. This laser generates laser pulses of approximately 70 fs with a repetition rate of 80 MHz. The wavelength is tunable between 720 nm and 850 nm. We normally use 780 nm in this research. The laser beam is splitted using a beam-splitter. This beam splitter transmits about 95% of the incident laser beam, and reflects the remaining 5%. The transmitted high power beam is used to demagnetize the sample as explained in section 3.2.2 and is called the pump-beam. The low power reflected beam is used to probe the magnetization using the MOKE and is called the probe-beam.

Because we are interested in the time dependent behavior of the magnetization, we
want to be able to control the time delay between the pump-beam and the probe-beam. This is done by a computer controlled delay line which consists of a retro reflector with an adjustable position. The retro reflector reflects the incoming beam, but shifts it over a small distance. In this way the incoming and reflected beam are parallel but shifted. This means that the position of the laser spot on the sample will not change when we the delay line is moved. The second delay line is not computer controlled and set at a fixed position during the measurements. The pump-beam is transported through two polarizers, which main purpose is to control the pump-power. The probe-beam is polarized and then directed through a photo elastic modulator (PEM) which purpose will be discussed in a moment. Finally, both beams are focused on the sample with a high aperture laser objective (HALO). The spot-size on the sample is approximately 20 µm. Because we use out-of-plane magnetized samples, the probe beam is directed perpendicular onto the sample, because in this way the sensibility of the MOKE is maximized for the out-of-plane component of the magnetization. The reflected pump-beam is blocked, while the reflected probe beam is directed towards the detector. By determining the Kerr rotation of the probe-beam the trajectory of the z component of the magnetization can be reconstructed. To get the best possible signal to noise ratio we use a double-modulation scheme that will be explained below.

The first modulation is introduced by the PEM. The PEM consists of a birefringent crystal that is periodically compressed in one dimension by a piezoelectric crystal. Due to this compression the optical path of one polarization component is slightly changed. This induces a periodical phase difference (retardation) between the polarization components of the laser pulse. The interaction of the different components in our set-up with the probe beam can be described using the Jones formalism introduced before. In this way the dependency of the detector signal on the magnetization can be calculated. Before entering the PEM, the probe beam is first linearly polarized along a direction that makes an angle of 45° with the modulation axis of the PEM. The Jones vector of the laser beam after the polarizer becomes: \( \frac{1}{\sqrt{2}} (1, 1)^T \). The Jones matrix of the PEM is given by:

\[
M_{PEM} = \begin{pmatrix}
1 & 0 \\
0 & e^{A_0 \cos(\Omega t)}
\end{pmatrix},
\]

in which \( \Omega \) is the modulation frequency and \( A_0 \) the maximum retardation. The outgoing beam is thus described by \( \frac{1}{\sqrt{2}} (1, e^{A_0 \cos(\Omega t)})^T \). \( A_0 \) is set to \( \pi/2 \), so the outgoing beam oscillates between right- and left-handed circular polarization.

Subsequently the beam is focused on the sample. The effect of the ferromagnetic sample can also be represented by a Jones Matrix:

\[
M_{FM} = \begin{pmatrix}
r_s & r_{sp} \\
r_{ps} & r_p
\end{pmatrix},
\]

in which \( r_s \) and \( r_p \) are the reflection coefficients for s- and p polarized light. The off-diagonal terms are proportional to \( \epsilon_{xy} \) which was discussed before, depends on
3.2. MEASUREMENT METHODS

the magnetization. Finally, the laser pulse reaches the detector where a voltage is produced which contains the information of the polarization of the pulse. The voltage given by the detector can be written as [40]:

$$V \sim \frac{\alpha E^2}{4} (|r_s|^2 + |r_p|^2) \{1 + 2\varepsilon \sin[A \cdot \cos(\Omega t)] + \sin 2(\theta + \Phi) \cos[A \cdot \cos(\Omega t)]\}, \quad (3.6)$$

in which $E$ is the intensity of the laser pulse, $\alpha$ describes the detector's sensitivity, $\varepsilon$ is the already introduced Kerr ellipticity, $\theta$ is the Kerr rotation, and $\Phi$ is the angle between the polarizer of the analyzer and the incident lights polarization axis. $\Omega$ is the modulation frequency of the PEM. Using Bessel functions, the terms $\sin[A \cdot \cos(\Omega t)]$ and $\cos[A \cdot \cos(\Omega t)]$ can be expanded in a harmonic series. In this way the voltage can be rewritten up to second order as:

$$V = V_0[1 + J_0(A) \sin(2(\theta + \Phi)) + 4J_1(A)\varepsilon \cos(\Omega t) + 2J_2(A) \sin(2(\theta + \Phi) \cos(2\Omega t) + ...], \quad (3.7)$$

with $J_n$ the nth order Bessel function and $V_0 = \frac{\alpha E^2}{4} (|r_s|^2 + |r_p|^2)$. We can rewrite this equation by collecting terms containing zeroth, first and second order harmonics in $\Omega t$, and get:

$$V = V_{static} + V_{1f} + V_{2f}, \quad (3.8)$$

with:

$$V_{static} = V_0[1 + J_0(A) \sin(2(\theta + \Phi))], \quad (3.9)$$

$$V_{1f} = V_04J_1(A)\varepsilon \cos(\Omega t V), \quad (3.10)$$

$$V_{2f} = V_02J_2(A) \sin 2(\theta + \Phi) \cos(2\Omega t). \quad (3.11)$$

The signal from the detector is analyzed by a first lock-in amplifier (L1). As a reference signal for L1 the 50 kHz signal from the PEM is used. The input voltage from the detector contains a 50 kHz and a 100 kHz component (respectively $V_{1f}$ and $V_{2f}$). So by selecting the first or second harmonic at the first lock in amplifier, we can measure $V_{1f}$ and $V_{2f}$ respectively, which are proportional to the Kerr ellipticity and the Kerr rotation and thus contains information about the magnetization.

To improve the signal-to-noise ration, a second modulation is used. The pump beam is chopped by the chopper at a frequency of 60 Hz. The output of the first lock in is then used as input for a second lock-in amplifier (L2). The frequency of the chopper is used as reference signal for L2. The L2 signal still contains the information from L1 about the Kerr rotation/ellipticy, but its signal to noise ratio is improved a lot. The main reason is that the second lock-in amplifier cancels drift that is present in the system due to for example instabilities in the optical components [8].
We put a lot of effort in optimizing the signal to noise ratio in measurements. One source of noise are reflections of the pump-beam into the detector. Therefore we made sure that we blocked all reflections of the pump-beam as much as possible. We also minimized the amount of reflections from the pump-beam that hit the detector using the polarizers of the pump-beam. By polarizing the pump-beam 90 degrees with respect to the polarizer in front of the detector we can minimize the amount of pump-beam reaching the detector. We also optimized the laser performance by calibrating it. The optical set-up is also re-calibrated and all the optical components are cleaned.

3.2.4 Kerr-Microscopy

Like the TR-MOKE set-up, the Kerr Microscope also uses the MOKE. It works basically as a normal microscope, but with added polarizers which allows for the detection of the MOKE. The microscope is equipped with a high intensity Xe arc, placed directly behind the aperture diaphragm. This diaphragm can be positioned in different positions, allowing to change the angle of incidence on the sample, and thereby changing the sensibility to the different types of MOKE. The advantage of the Kerr-Microscope is that it gives a spatial overview of the magnetization. The CCD camera employed in the Evico Kerr microscope used allows us to record images at 16 frames per second with a resolution of 100 nm, allowing to study dynamics processes on the second-timescale. The Kerr-Microscope cannot be used to investigate the fast magnetization processes, but it can be used to characterize the magnetic switching behavior, as discussed in chapter 4. An example of a image from the Kerr microscope is given in figure 3.9. To improve the measurements, first a back ground image is recorded. During the measurements, this background image is subtracted from the recorded images. This leads to an improved sensitivity to changes in the magnetization.
Figure 3.9 – Example of an image recorded by the Kerr microscope. Depicted is a nano-wire connected to a contact, which is magnetic too. The dark area indicates regions with magnetization in the positive z direction, the light area means magnetization in the opposite direction.
Chapter 4

Current assisted switching measurements

In chapter 1 we discussed current assisted magnetization reversal measurements performed by Miron et al on Pt/Co/AlOx samples [16]. They observed a field-like effect of the current on the magnetization reversal. With 'field-like' we mean that the current wants to stabilize one magnetization direction just like an applied magnetic field. Miron proposed two possible effects, both due to the presence of spin-orbit interaction, that have the right symmetry to be able to explain their measurements: the spin Hall effect (SHE) and a Rashba interaction. We discussed both spin-orbit (s-o) torques in chapter 2. There we argued that the Rashba field should not be present in systems with two identical interfaces like Pt/Co/Pt and Pt/CoB/Pt. This was confirmed by Haazen et al [17] in this group for Pt/Co/Pt nano-wires. Also in this group, Murè et al discovered an effect with the same symmetry as the effect found by Miron in de-pinning experiments on Pt/Co/Pt nano-wires.

We performed measurements to verify the presence of this s-o torque in Pt/CoB/Pt and Pt/CoB/AlOx nano-wires. In the first part of this chapter we discuss performed measurements in which we switch the magnetization of our nano-wires with an external magnetic field, and measure the change in switching field due to an applied current. We saw in chapter 2 that according to macro-spin simulations the torque induced by the SHE should be able to assist or oppose this field-induced switching. In chapter 2 we explained that in a Pt/CoB/Pt wire the torque exerted by the spin-current originating from the two Pt layers oppose each other. To maximize the torque exerted on the ferromagnetic layer, we first performed measurements on Pt/CoB/AlOx wires. However, we will see that we don’t observe a measurable change in switching field with the symmetry of the SHE. Subsequently we performed measurements on Pt/CoB/Pt wires in which we eliminate the Rashba field. We compared those measurements with simulations, to find an estimate for the Hall angle in our system.

In the second part of this chapter, we tried to switch the magnetization using only
the SHE and an in-plane applied magnetic field. This leads to interesting observations as we will show in section 4.2.

The current we send through the nano-wires will generate an Oersted field and Joule heating, which can influence the domain wall dynamics. These effects can easily be separated from the SHE, because they will be symmetric with respect to the current direction. We will take a closer look at those phenomena in section 4.3.

4.1 Field-induced current-assisted switching measurements

In chapter 2 we explained that, according to macro-spin simulations, the torque exerted by the injected spins originating from the SHE could help or oppose field-induced magnetic switching when an in-plane magnetic field in the x direction (in the current direction, see figure 4.1 b) is applied. This means that measuring the change in switching field as a function of applied current could provide us with information about the presence of the SHE in our systems. Like in the performed simulations, we switch the magnetization of an out-of-plane magnetized sample using a magnetic field that is tilted by a certain angle $\Theta$ with respect to the out-of-plane axis (z axis) of the sample, see figure 4.1b. This means that the applied field can be decomposed in an out-of-plane component, $H_z$, and an in-plane field component. This in-plane component can lie either along the x direction, which means that the field lies in the x-z plane ($H_{xz}$ field) or along the y direction, which means that the field lies in the y-z plane ($H_{yz}$ field) as indicated in figure 4.1b.

In chapter 2 we also discussed that the effect of the SHE on the magnetic switching changes sign with inverting the current direction or the in-plane field direction. Those symmetry properties of the SHE provide us with a good test to check whether any effect we measure on the switching field can be explained by the SHE.

We investigated magnetic reversal in three different systems: Pt(5)/CoB(0.8)/AlOx(0.89), Pt(5)/CoB(0.7)/Pt(2) and Pt(5)/CoB(0.7)/Pt(5). We will start this section to show results performed on the first system with an AlOx toplayer. Because in this system only the bottom layer is made of platinum and the top layer is isolating, there will be only an injected spin-current originating from the bottom layer, and no spin-current with opposite polarization from the top-layer. So we expect a maximal torque exerted by the SHE in this sample. However, because this system contains two unequal boundaries, we could also expect a stronger Rashba field, which has the same symmetry with respect to the current direction as the SHE. This means that in this sample it could be hard to distinguish between the two spin-orbit torques. The second system we will discuss contains a second Pt layer replacing the AlOx having a thickness different from the bottom Pt layer. Because we now have two equal interfaces we cancel the Rashba field and we will be able to study the contribution from the SHE. The last sample
contains two equal platinum layers. Now we expect the two spin currents originating from the two platinum layers to be equal and to cancel out each other’s torque on the magnetization. So in this sample we expect no change in switching field due to the SHE.

Before we discuss the measurement results, we will first explain how we can interpret the data obtained from switching field measurements.

4.1.1 Decomposing the change in switching field

The blue curve marked with (blue) squares in figure 4.1a shows a typical hysteresis loop for a ferromagnetic material. We will now discuss how this loop changes when we apply a current, and how we can extract information about the effect of s-o torques on the magnetic switching from the altered hysteresis loop. When we apply a (positive) current we will get the red curve marked with red triangles: the magnetization is switched at a lower field. We define this change in switching field4.1b as $\Delta H^{Total}$, which will have a negative value when the magnetization is switched at a lower field due to the current. This total change in switching field can be written as the sum of a symmetric part,$\Delta H^{Sym}$, which does not change when the current is inverted, and an asymmetric part,$\Delta H^{Asym}$, which changes sign when the current is inverted (see figure 4.1a):

$$\Delta H^{Total} = \Delta H^{Sym} + \Delta H^{Asym}. \quad (4.1)$$

In chapter 2 we saw that the effect of the s-o torques changes sign when the current direction is inverted. So the s-o torques will have a major contribution to $\Delta H^{Asym}$. For Pt/CoB/Pt, we will ascribe $\Delta H^{Asym}$ completely to the SHE, because we can neglect the Rashba interaction. The symmetric part, $\Delta H^{Sym}$, is due to Joule heating and Oersted fields. Heating up the sample will lower the total magnetization and therefore the magnetization is switched more easily. Besides that, thermal fluctuations will accelerate domain nucleation. We will come back to Joule heating in more detail later on in section 4.3, were we will also explain why the effect of the Oersted field is symmetric with respect to the current direction.

We are mainly interested in $\Delta H^{Asym}$, so we need a method to extract its value from $\Delta H^{Total}$. The solution is to measure the change in switching field for both positive and negative current. Because the symmetric contribution changes sign when we invert the current direction, we can write:

$$\Delta H^{Asym}(I > 0) = -\Delta H^{Asym}(I < 0). \quad (4.2)$$

The contribution of the symmetric effect is unaffected by the direction of the current:

$$\Delta H^{Sym}(I > 0) = \Delta H^{Sym}(I < 0). \quad (4.3)$$
CHAPTER 4. CURRENT ASSISTED SWITCHING MEASUREMENTS

Figure 4.1 – a) A typical switching measurement without current and with a positive applied current. Indicated are symmetric and asymmetric part of the change in switching field. b) Definitions of an applied $H_{xz}$ and an applied $H_{yz}$ field. The angle of the applied magnetic field with the z axis is indicated with $\Theta$.

So if we measure the change in switching field for positive and negative current and subtract both results we cancel the symmetric part:

$$\Delta H^{Total}(I > 0) - \Delta H^{Total}(I < 0) = 2\Delta H^{Asym}(I > 0).$$  \hspace{1cm} (4.4)$$

If we add them we will cancel the asymmetric part:

$$\Delta H^{Total}(I > 0) + \Delta H^{Total}(I < 0) = 2\Delta H^{Sym}(I > 0).$$  \hspace{1cm} (4.5)$$

In this way we can separate the symmetric and asymmetric change in switching field, which gives us the opportunity to isolate the effect from the s-o torque on the switching field.

In the next subsections we will plot $\Delta H^{Sym}$ and $\Delta H^{Asym}$ in separate graphs. We will plot $\Delta H$ as a function of a positive current, for an in-plane applied magnetic field in the positive x or y direction, unless stated otherwise. Remember that we always use the system of coordinates depicted in figure 4.1b.
4.1. FIELD-INDUCED CURRENT-ASSISTED SWITCHING MEASUREMENTS

![Graphs](image)

**Figure 4.2** — a) Change in switching field measured on a 50 µm wide Pt(5)/CoB(0.7)/AlOx(0.89) nano-wire. Plotted is the asymmetric part of the change in switching field, for or both a \( H_{xz} \) and a \( H_{yz} \) applied magnetic field. Measurements done with a magnetic field making an angle of 80° with the z axis. b) Symmetric part of the change in switching field due to Joule heating and Oersted fields.

### 4.1.2 Pt(5)/CoB(0.8)/AlOx(0.89)

The first measurements we show are done on a 50 µm wide Pt(5)/CoB(0.8)/AlOx(0.89) nano-wire. As mentioned before, due to the absence of a second platinum layer we expect to maximize the total injected spin current. So we expect a maximum effect on the switching field in this sample. We performed measurements with an applied \( H_{xz} \) field as well as measurements with an applied \( H_{yz} \) field. The angle \( \vartheta \) of the applied field with the z axis was 75°.

The results are shown in figure 4.2. In figure 4.2a we plotted \( \Delta H_{\text{Asym}} \) and in figure 4.2b we plotted \( \Delta H_{\text{Sym}} \), both as a function of the current density. We have plotted results for measurements done with \( H_{xz} \) fields as well as from measurements done with \( H_{yz} \) fields. We see a clear symmetric contribution to the total change in switching field as a function of the current density in figure 4.2b. At a current density of \( 6 \cdot 10^{10} \text{A/m}^2 \) we observe a \( \Delta H_{\text{Sym}} \) of -7 mT. As expected, this contribution does not depend on the direction of the in-plane field. However, we observe almost no asymmetric effect, as becomes clear from figure 4.2a. Possibly, the Rashba field present in this system because of the asymmetric boundaries of the magnetic system is responsible for this unexpected behavior, but at this point we cannot explain those observations. We will now take a look at Pt/CoB/Pt samples, were we exclude the Rashba field.
CHAPTER 4. CURRENT ASSISTED SWITCHING MEASUREMENTS

Figure 4.3 – a) Change in switching field in 20 µm wide Pt(5)/CoB(0.7)/Pt(2) nanowires. Plotted is the change in switching field due to the SHE, for both a $H_{xz}$ and $H_{yz}$ applied magnetic field. b) The symmetric part of the change in switching field, due to Joule heating and Oersted fields.

4.1.3 Pt(5)/CoB(0.7)/Pt(2) wires

As we saw in the previous subsection, we did not detect the SHE in the Pt(5)/CoB(0.8)/AlOx(0.89) nano-wire. We will now show measurements done on a 20 µm wide Pt(5)/CoB(0.7)/Pt(2) nano-wire. Due to the difference in thickness of the two Pt layers, we expect a net spin current into the magnetic layer induced by the SHE. In figure 4.3 we see the result of switching field measurements done on this sample. We performed the same measurements as on the Pt/CoB/AlOx sample, but now using an angle of 80° of the magnetic field with the $z$ axis. Again, we performed measurements with an applied $H_{xz}$ field as well as an applied $H_{yz}$ field. In figure 4.3 b we see the symmetric part of the change in switching field and in figure 4.3 a the asymmetric part. Again we see a symmetric contribution to the switching field that is independent of the direction of the in-plane magnetic field. But now we also see a more interesting asymmetric contribution appearing. For the $H_{xz}$ applied field, we see a clear asymmetric contribution to the change in switching field, which increases with increased current density. The size of this change in switching field is about -3.5 mT at a current density of $4 \times 10^{10}$ A/m². The switching field without applied current is approximately 17 mT. Remember that this field is the total field, which has both an in-plane and an out-of-plane component. For the $H_{yz}$ field we also see an asymmetric change in switching field, but the effect is much smaller. Because simulations showed no effect on the switching field in the $H_{yz}$ field configuration (see chapter 2), the measured effect is probably due to a misalignment of the applied magnetic field.
4.1. FIELD-INDUCED CURRENT-ASSISTED SWITCHING MEASUREMENTS

Figure 4.4 – a) Change in switching field as a function of current, for both positive and negative applied x field. b) Hysteresis curves measured in the Kerr microscope with only an out-of-plane external magnetic field. (27 mA is equal to $1.75 \cdot 10^{11}$ A/m in this sample, a 20 μm wide Pt(5)/CoB(0.7)/Pt(2) nano-wire ).

We observed a current induced effect that inverts its direction when the current direction is inverted, but the effect should also change sign with inverted in-plane field direction. In figure 4.4 we see $\Delta H_{\text{Asym}}$ as a function of the applied current for an applied field along the positive and the negative x direction. We see that the effect is indeed inverted. In figure 4.4b we also show a hysteresis curve measured on the Kerr microscope. In this set-up it is possible to apply a pure out-of-plane magnetic field. We see that with no in-plane field, positive and negative current have the same effect on the switching: In this case we only have the effect of the Oersted field and Joule heating, the symmetric contributions. So indeed we need an in-plane field in the x direction to have an effect of the SHE on magnetic switching.

From those observations, we conclude that the SHE can indeed explain the current assisted magnetic switching measurements in Pt/CoB/Pt nano-wires. In the last part of this section, we will show that the relative thickness of the two platinum layers is indeed important for the size of the induced torque as argued in chapter 2.

4.1.4 Pt(5)/CoB(0.7)/Pt(5).

If we fabricate a sample with two platinum layers having the same thickness we expect the net torque from the SHE to be zero. In figure 4.5a we show the results of measurements on a Pt(5)/CoB(0.7)/Pt(5) wire. We used an angle of the magnetic field with the z axis of 80°. We see that the effect is indeed small, but still approximately 1 mT at a current density of $6 \cdot 10^{10}$ A/m². This is probably due to a difference in
growth between the two layers, which could affect the Hall angle of the material. It could also be a non-zero Rashba field. In figure 4.5b also the symmetric part of the results is shown. These results are roughly the same as for the Pt(5)/CoB(0.7)/Pt(2) wire. The difference is explained in section 4.3.

To summarize, those measurements showed that it should be possible to control the net spin Hall torque by varying the platinum thickness. We cannot unlimited make the Hall angle larger by increasing the layer thickness. This process is limited by spin accumulation as described in chapter 2.

### 4.1.5 Estimate of the effective Hall angle

In chapter 2 we showed results of macro-spin simulations of current assisted magnetic switching. Those simulations were done for Pt/CoB/Pt nano-wires. A Hall angle for both Pt layers was determined depending on their thickness. Using the Hall angle of both layers, a net spin current injected into the magnetic layer was calculated. We also discussed that we performed simulations for different values of the Hall angle. We started with the values found in literature, but also performed simulations with higher and lower values. The results were given in figure 2.14.

We will now compare those simulations with the measurements we performed in this section on the Pt(5)/CoB(0.7)/Pt(2) nano-wire to determine the effective Hall angle of our system. As we saw in figure 2.14, the simulations show a linear dependence of
4.2 CURRENT INDUCED MAGNETIC SWITCHING

In the previous section we saw that the SHE in combination with an applied magnetic field in the x direction could assist or oppose field induced magnetic switching. In this section we investigate whether it is possible to switch the magnetization using only an applied current in combination with a pure in-plane magnetic field to switch the magnetization of a perpendicularly magnetized nano-wire (Pt(5)/CoB(0.7)/Pt(2) in...
Before we discuss the performed measurements, we will show that when a large pure in-plane magnetic field is applied, the system will be in a multi-domain state. This has some important consequences.

4.2.1 Multi-domain state

In order to perform good quality switching experiments, we need to guarantee that the system is in a well-defined single-domain state. When we apply large, pure in-plane magnetic fields, we face one complication. The magnetization will break up in a multi-domain state. Figure 4.7b shows a Kerr image of a part of the wire with a large applied in-plane magnetic field. We see that the system is indeed in a multi-domain state and that the domains are very small compared to the laser spot size (indicated by the circle). After we remove the in-plane magnetic field, the system remains in this multi-domain state. In the TR-MOKE set-up we measure the average magnetization over the whole laser spot. This average magnetization will be zero in the multi-domain state. This means that after we apply a large in-plane magnetic field the measured remanence (magnetization at zero field) will be zero. So we have to take care when we perform measurements with large in-plane magnetic fields, because the multi-domain state is not as well defined as a single domain state.

We can use the fact that the remanence is zero after applying a large pure in-plane field to tilt our magnet exactly in-plane. In figure 4.7a we see the procedure to do so in the TR-MOKE setup. We make a field-sweep and look at the remanence. If there is a small out-of-plane component of the magnetic field, the system prefers magnetic domains which magnetization points along this out-of-plane component and we measure a non-zero remanence. We keep decreasing the angle of the applied magnetic field until the remanence is zero. At that point we are sure that the field is completely in plane.

4.2.2 Current induced switching measurements

We will now show two types of experiments using an applied in-plane field in combination with an applied current. During both experiments, we performed field-sweeps. Those field-sweeps can be forward, meaning going from low to high magnetic field, or backwards, going from high to low magnetic field. In the first experiments, we prepare the system in a multi-domain state, and look at the effect of an applied current on this state. We expect to see a field-like effect from the SHE which would give the magnetization a preferred direction and bring it into a single-domain state. The second experiment is the more interesting one. Here we will bring the system in a single-domain state, and try to switch its magnetization with only an applied current
in combination with an in-plane magnetic field. Note that the magnetic field in all the measurements performed in this section is directed along the x axis.

Field like effect on the multi-domain state

In the first experiment, we prepare the system in the multi-domain state by applying a large in-plane magnetic field. We then perform a backward field-sweep were we slowly decrease the field to zero and increase it again in the opposite direction. The field is purely in-plane during the whole measurement. The results are given in figure 4.8 for different values of the applied current. Note that in this section we used total current instead of current density. To compare the results in this section with results from section 4.1, we give the current densities for the applied currents in this section in table 4.1. When no current is applied, the system stays in the multi-domain state during the whole field-sweep. This means we measure zero magnetization during the whole measurement. The slope of the curve is due to Faraday [38] rotation. However, when we apply a current, the torque induced by the SHE in combination with the in-plane magnetic field tends to tilt the magnetization in a preferred direction, in this case along the positive z axis. At large in-plane fields, this torque cannot compete with the in-plane field which keeps the system in the multi-domain state, but when the field is decreased, we observe a net out-of-plane magnetization. This means that the
domains with a magnetization in the positive z direction start to grow. The remanence at zero field is not longer zero. The measured remanence increases with increasing current, but even at the largest applied current, 20 mA, the saturation magnetization is not reached, which means that there still exist domains with opposite magnetization. Unfortunately we could not increase the current any further due to Joule heating.

In the previous section, we saw that the current induced effect changes sign when the in-plane component of the magnetic field is switched. We also observe this behavior in figure 4.10. When the in-plane field crosses zero and is increased in the opposite direction, we see that the magnetization direction switches at some point. However, we don’t observe a nice switch, because the magnetic field starts to bring the system in a multi-domain state again.

In figure 4.9 we show results from the same kind of measurements. Here we show both a forward and a backward scan, for both positive and negative applied current. These measurements clearly show that the effect is inverted with inverted in-plane magnetic field, but also with inverted current direction.

**SHE induced switching from a single-domain state**

The previous measurements clearly show that the current induces a field-like effect on the magnetization. However, those measurements are performed on a multi-domain state which is not very well defined. Now we will look whether it is possible to switch a well defined single-domain state using a current and an in-plane magnetic field. We prepare the system in a nice single-domain state by tilting the applied field out-of-plane. Then we switch off the field and turn the magnet in plane and perform a forward or backward field-sweep starting from zero field. We search for a combination of in-plane field and current for which we can switch the magnetization. The results are shown in figure 4.10. The blue circular curve shows the magnetization as a function of in-plane field when no current is applied. We see that we start from a single domain state at zero field, for which we measure the saturation magnetization. For large fields we end up in a multi-domain state. When we apply a negative current and perform a forward field-sweep, we observe that for a large enough in-plane field the magnetization switches. Again we observe that the direction of the effect depends on the current and in-plane field direction. Unfortunately the in-plane field for which the switching occurs is close to the field for which we start creating the multi-domain state. This means that the system is not in a well defined state when we switch the magnetization. We could increase the current to switch at lower fields, but due to Joule heating at such high currents we damage the samples.
4.2. CURRENT INDUCED MAGNETIC SWITCHING

Figure 4.8 – Backward in-plane field scan for different applied currents. The in-plane field is directed along the x axis.

<table>
<thead>
<tr>
<th>Current (mA)</th>
<th>Current density (A/m²)</th>
</tr>
</thead>
<tbody>
<tr>
<td>20</td>
<td>1.3 \cdot 10^{11}</td>
</tr>
<tr>
<td>10</td>
<td>6.5 \cdot 10^{10}</td>
</tr>
<tr>
<td>5</td>
<td>3.3 \cdot 10^{10}</td>
</tr>
</tbody>
</table>

Table 4.1 – Table with the current densities for the used currents in this section.

Figure 4.9 – a) Backward field scan for positive and negative current. b) Forward field scan for positive and negative current. The in-plane field is directed along the x axis in both graphs.
CHAPTER 4. CURRENT ASSISTED SWITCHING MEASUREMENTS

Figure 4.10 – Field sweep with an in-plane applied field for 0 mA, 20 mA and -20 mA. The in-plane field is directed along the x axis.

4.3 Contributions to the symmetric part of the switching field

In section 4.1 we saw that besides the effect from the s-o torque on the switching field, we also observe an effect which is symmetric with respect to the current direction. In this section we will take a closer look to the mechanisms behind this effect. We will see that there are two contributions in our nano-wires which can explain the effect: Joule Heating and Oersted fields.

4.3.1 Joule heating

We will first take a look at Joule heating. From the Bloch law we know that the magnetization of a ferromagnet reduces by increasing the temperature.

\[
\frac{M(T)}{M(0)} = 1 - \left( \frac{T}{T_c} \right)^{3/2},
\]

in which \(M(T)\) is the magnetization at temperature \(T\), \(M(0)\) is the magnetization at zero temperature, and \(T_c\) is the Curie temperature. This will also affect the shape anisotropy discusses in chapter 2, and thereby the switching field. Figure 4.14b shows the result of a macro-spin simulation we performed of the change in switching field as a function of temperature for a Co magnetic layer with out-of-plane anisotropy. Those simulations only take the decreased magnetization due to an increased temperature according to equation 4.6 into account, and neglect totally domain nucleation by thermal fluctuations, which is a very important mechanism in magnetization reversal. Therefore those simulations are only a very rough estimation. We indeed see a mea-
4.3 CONTRIBUTIONS TO THE SYMMETRIC PART OF THE SWITCHING FIELD

Figure 4.11 – Plot of $H_z$ vs $H_x$. Increasing the in-plane component of the magnetic field decreases the out-of-plane field needed to switch the magnetization. Measured on a Pt(5)/CoB(0.7)/Pt(2) nano-wire.

Measurable change in the switching field already for not too high temperatures. To confirm that Joule heating really affects the magnetic switching, we have to get an estimate of the temperature in the magnetic wires. We performed both simulations with Comsol multiphysics and measurements on a Pt/CoB/AlOx sample. More details about the simulations and temperature measurements are given in Appendix A. The results are plotted in figure 4.14a. We have plotted the increase of temperature as a function of applied current density. We measured a maximum increase in temperature of 20 K for the currents used in this research. Comparing this with the switching simulation we conclude that we can expect changes in the switching field in the order of magnitude of 1 mT due to Joule heating.

4.3.2 Oersted fields

The second mechanism that can cause a change in switching field which is symmetric with respect to the current direction is the generation of Oersted fields inside the magnetic wire by the electric current. We performed simulation with Comsol to investigate the shape of the fields inside the wire. The results are shown in figure 4.12. The figure only shows the edges of the wire. We see that the magnetic field is in-plane along the y direction inside the magnetic layer, except at the edges where there is an out-of-plane component. This in-plane magnetic field can play a crucial role in the switching of the magnetization. The size of the out-of-plane component needed to switch the magnetization depends on the in-plane component of the field. To prove this, we measured the switching field for different angles on a Pt(5)/CoB(0.7)/Pt(2) nano-wire and plotted the out-of-plane switching field versus the in-plane field in figure 4.11. So if Oersted fields contribute to the in-plane component of the magnetic field, we would need a smaller out-of-plane component to switch the magnetization.
If we take a closer look at the symmetric part in the change of switching field discussed in section 4.1, we see that the effect is the largest for the Pt(5)/CoB(0.8)/AlOx(0.89) (-6 mT at $5 \cdot 10^{10} \text{A/m}^2$) sample, and smallest for the Pt(5)/CoB(0.7)/Pt(5) (-2 mT at $5 \cdot 10^{10} \text{A/m}^2$) sample. For the Pt(5)/CoB(0.7)/Pt(2) wire we find an intermediate value of -4 mT at $5 \cdot 10^{10} \text{A/m}^2$. We can explain this using figure 4.12. For the Pt(5)/CoB(0.7)/Pt(5) wire, the magnetic layer is in the center, where the field is almost zero. For the Pt(5)/CoB(0.7)/Pt(2) sample and even more for the Pt(5)/CoB(0.8)/AlOx(0.89) sample the magnetic layer is more towards the edge, where the in-plane magnetic field is larger. In the Pt(5)/CoB(0.8)/AlOx(0.89) sample we can also expect a additional Rashba field, which is in the y direction, which could also contribute to the in-plane magnetic field.

The effect of the out-of-plane component of the Oersted field on the edges of the sample on the magnetic switching can be shown with images taken at the Kerr microscope and from simple simulations. Two of the Kerr images are shown in figure 4.13. We see the 50 µm wide wire, and on the right side a part of the larger contacts which are of the same composition as the wire itself (see chapter 3). In the first image, we see switching without any current. The switching starts with nucleation in the contacts and continues by domain-wall motion until the complete wire is switched. Because the switching is due to nucleation at lower anisotropy parts and the followed domain-wall motion, the measured switching field is lower than values found by (macro-spin) simulations. If we now send a current trough the wire, the switching behavior looks like depicted in figure 4.13b. The out-of-plane component of the Oersted field helps the switching in one half of the wire and oppose it on the other half. The rest of the wire is again switched by domain wall motion. This switching occurs before the nucleation observed at the contacts in the no current case, so the out-of plane component of the Oersted field clearly decreases the required switching field. If the current direction is reversed, the other half of the wire is switched first, but the switching is assisted in both situations and the effective change in switching field doesn’t depend on the sign of the current.

We can conclude that both the Joule heating and the Oersted field have an measurable effect on the switching field. In the Pt(5)/CoB(0.8)/AlOx(0.8) sample we can also expect an effect of the Rashba field. These effects are however pretty complex, and therefore it is hard to predict the size of each individual contributions on the magnetic switching.
4.3. CONTRIBUTIONS TO THE SYMMETRIC PART OF THE SWITCHING FIELD

Figure 4.12 – Simulations of the magnetic field inside the wire. We see a cross section of the wire, and we plotted only the two edges. The arrows show the field direction, and the field strength is represented by both the color and the size of the arrows. The direction of the current is into the paper. The layer thicknesses are indicated in nm.

Figure 4.13 – Kerr images of the magnetic switching in Pt/CoB/Pt nano-wires. We see the nano-wire and a small part of the contact. In a) switching with only a magnetic field is depicted. In b) the magnetic switching is current assisted. More details are given in the main text.
Figure 4.14 – a) Simulations and measurements of the temperature in a Pt(5)/CoB(0.7)/AlOx(0.89) wire b) Simulations of the effect of the temperature on the switching field. More details about the temperature measurements and simulations are given in appendix A.
Chapter 5

Current influenced magnetic damping

In the previous chapter we characterized the SHE in Pt/CoB/Pt nano-wires by measuring current assisted magnetic switching. In chapter 2 we argued that the torque induced by the SHE could also influence the magnetic damping. This could provide us with a new measurement technique to study current induced spin-orbit torques. Therefore we performed damping measurements which are discussed in this chapter.

The first measurements in this chapter are performed without the injection of electrical currents, to characterize our samples. In chapter 2 we saw that the damping in thin magnetic films can be enhanced by spin pumping. We argued that angular momentum is transported from the magnetic layer to an adjacent non-magnetic layer. This non-magnetic layer, Pt in our case, acts as a spin sink for the injected angular momentum. We want to check whether the Pt layer thickness influenced the ability of the Pt to absorb the injected spins. We did this by measuring the damping as a function of the top and bottom Pt layer thickness. Those measurements are performed on films instead of nano-wires to optimize the signal to noise ratio.

The final measurements are performed on magnetic nano-wires to study whether we can affect the damping in a measurable way using the SHE. In chapter 2 we argued that the spin current injected by the SHE could exert a torque on the magnetization that is able to influence the magnetic damping, when a magnetic field with a component in the y direction is applied. This was confirmed by macro-spin simulations. Therefore we performed damping measurements on the same nano-wires we used in chapter 4. We performed those measurement on both Pt/CoB/Pt and Pt/CoB/AlOx nano-wires.

In the first section, we will discuss the field dependency of the damping parameter. In section 5.2 the influence of the power of the pump-beam on the measurements is discussed. In the third section we discuss the spin sink study on thin films and in the last section we will discuss the current assisted magnetic damping measurements.
5.1 Field depended damping measurements

In this section we discuss the field depended damping measurements we performed in this chapter. In chapter 3 we discussed that we need to use an external magnetic field to tilt the magnetization partly in-plane to be able to perform damping measurements. It turns out that $\alpha$ depends on the size of the applied magnetic field. This behavior is not described by the macro-spin model, but by a so called multi-spin model introduced by Waloski et al [41]. We will now take a closer look to this model.

In figure 5.1d, a typical measurement of $\alpha$ as a function of applied field is showed. Those measurements are performed on a Pt(5)/CoB(0.7)/Pt(5) thin film using a field applied with an angle of 82.5 degrees with the z axis (see figure 5.1c). At low magnetic fields we observe an enhanced $\alpha$. This field depended behavior of the damping parameter is described by the multi-spin model. It turns out that the anisotropy is not constant throughout our samples. The sample is divided into small domains, each with their own anisotropy, see figure 5.1a. The anisotropy varies typically a few percent around the average value [8]. As we will see in a moment, the field dependency of the damping is a result of those anisotropy differences between the domains.

In chapter 2 we saw that the magnetization dynamics are described by the LLG equation, which contains an effective magnetic field $H_{eff}$. When the magnetization is out of equilibrium, the effective magnetic field exerts a torque which induces a precessional motion of the magnetization. Due to damping, the magnetization will relax back to equilibrium. The effective field of our samples is determined by the applied field and the anisotropy field, see figure 5.1c. As showed by Kuiper [8], the precession frequency of the magnetization around the effective magnetic field scales with the sum of the applied magnetic field and the anisotropy field:

$$\omega = \frac{\gamma \mu_0}{1 + \alpha^2} \left( H_{app} \frac{\cos(\beta)}{\cos(\theta)} + H_K - M_{sat} \right), \quad (5.1)$$

in which $\beta$ is the angle of the applied magnetic field with the z axis, $\Theta$ is the angle of the magnetization with the z axis (see figure 5.1c), $H_{app}$ is the applied magnetic field, $H_K$ is the anisotropy field and $M_{sat}$ is the saturation magnetization. Using equilibrium conditions, Kuiper found that the precession frequency can also be written as:

$$\omega = \frac{\gamma \mu_0}{1 + \alpha^2} \eta H_{app} \sin(\beta) \sin(\theta). \quad (5.2)$$

Because the equilibrium angle $\Theta$ depends on the anisotropy, from this equation we obtain an $\omega$ that depends on the anisotropy. This means that the individual domains with varying anisotropy will precess with their own frequency $\omega$. The measured signal is the sum over the damped oscillation of all individual domains covered by the laser spot, which is much larger than the domains for thin magnetic films, see figure 5.1a. This leads to an increased apparent damping resulting from the line broadening and different phase in frequency space [41].
Figure 5.1 – a) The sample is divided in 'domains' with slightly different values of the anisotropy. Adapted from [41]. b) The angle $\Theta$ the magnetization makes with the $z$ axis as a function of applied field, determined for Pt/CoFeB/Pt films by Kuiper [8]. c) Definitions of $\beta$ and $\Theta$. In equilibrium, $M$ lies along the effective field direction, which is a combination of the applied field and the anisotropy field. d) Typical measurement of $\alpha$ as a function of the applied field.
From equation 5.2 we know that $\omega$ depends on $\Theta$, the angle of the magnetization with the z axis. Figure 5.1b shows $\Theta$ as a function of the applied magnetic field. The angle $\Theta$ depends on both the applied magnetic field and the anisotropy field. In a sample with strong anisotropy, the field has to be large to tilt the magnetization in the direction of the applied field. When we apply a very large field, the angle of the magnetization will approach the angle of the applied field. At those large fields, variations in the anisotropy will have a smaller effect on $\Theta$ than at a small applied field. This means that at low field, the variation of $\Theta$ and thus $\omega$ between the domains is larger, and thus the damping parameter is distorted more at lower fields, which leads to the field-depend damping parameter showed in 5.3b.

5.2 Optimizing the pump beam power

As explained in chapter 3, in the TR-MOKE set-up a pump-beam is used to demagnetize the magnetic sample and bring the magnetization out of equilibrium. It turns out that we have to tune the pump-power carefully to obtain good measurements. Obviously, if we use a pump-power that is too low, we do not excite the system enough to be able to measure nice oscillations. However, when we use a pump-power that is too high, the quality of the signal will decrease as we will see in a moment. The optimal pump-power depends on the system we use. Therefore, we individually optimized the pump-power for our Pt/CoB/Pt as well as our Pt/CoB/AlOx samples.

In figure 5.2a we show damping measurements performed on a Pt(5)/CoB(0.7)/Pt(2) sample using three different values for the pump-power and an applied field under an angle of 82.5 degrees with the z axis. As we expected, when we use a pump-power that is too low, we do not excite the system enough to obtain a nice oscillation. However, when the pump-power is increased to much, the quality of the magnetic signal gets worse and we are not able to determine $\alpha$ accurately. The damping parameter as a function of applied field for a Pt(5)/CoB(0.7)/Pt(2) is plotted in figure 5.3a and b. We observe the decrease in damping for large magnetic field strengths that we discussed in the previous section. In figure 5.3a the results from measurements done with a pump-power of 25 mW are plotted, and in figure 5.3b from measurements done with a pump-power of 54 mW. The results are significantly improved when a pump-power of 54 mW is used. We decided to use 50 mW pump-power on the Pt/CoB/Pt samples. The same optimization is carried out for the Pt/CoB/AlOx samples. For those systems the signal is optimized using a pump-power of 25mW.

There could be a couple of reasons for the fact that we decrease the signal quality when we use a pump-power that is too high. Reflections of the pump-beam that hit the detector can introduce noise. Therefore, if we increase the pump-power we will also increase the noise. The second possibility is that we decrease the magnetization of the sample by increasing the temperature. To test this we measured a hysteresis
5.3 DAMPING MEASUREMENTS ON MAGNETIC THIN FILMS

Figure 5.2 – a) Measurement of the time evolution of the z component of the magnetization for different values of the pump-power. Measured on a Pt(5)/CoB(0.7)/Pt(2) sample. b) Hysteresis loop measured on a Pt(5)/CoB(0.8)/AlOx(0.89) nano-wire with low (20 mW) and high (71 mW) pump-power.

loop on a Pt/CoB/AlOx sample. On this sample we found that the signal quality decreases already at low pump-power compared with the Pt/CoB/Pt sample, and that the decrease in signal quality is also more drastic. We plotted the results in figure 5.2b. We observe a nice hysteresis loop for a pump-power of 20 mW. However, when the pump-power is increased to 71 mW, we observe a lot more noise, and a lower switching field. The lower switching field is due to the increased thermal fluctuations which enhances domain nucleation. However we do not observe a much lower magnetization, so we don’t expect this to be the reason for the decreased signal quality.

5.3 Damping measurements on magnetic thin films

In chapter 2 we discussed that an important extrinsic contribution to the magnetic damping in thin ferromagnetic layers originates from spin-pumping, as shown by Tserkovnyak et al [42], who calculated the damping parameter as a function of the ferromagnetic layer thickness. Spins are pumped into adjacent non-magnetic layers and therefore the increase in $\alpha$ depends on the ability to absorb spins of the adjacent layers. We wanted to investigate whether varying the Pt layer thickness, in the working range of this research, influences this spin-sink ability of the Pt layer. To study this we have performed measurements on samples with varying Pt layer thicknesses. We varied both the bottom and the top Pt layer thickness. The measurements were performed on two kind of samples. We will show that the choice of the substrate is
very important for the signal quality.

At first, we fabricated samples on Si wafers with only very thin, natural oxidation (10 nm). The reason to use this substrate is that SiOx has very low heat conductivity. Because in this study we are interested in current induced effects, Joule heating plays an important role as we discussed in the previous chapter. On top of the Joule heating, in the damping measurements we also use the high power pump-beam. Therefore we want to use a substrate with an oxide layer that is as thin as possible in order to optimize the heat flow from the wire through the substrate. We start with measurements we performed on thin films sputtered on those substrates. In figure 5.4 we have plotted the damping parameter as a function of the applied field. Those measurements are performed with an angle of 82.5 degrees with the z axis. We observe the expected field dependence of $\alpha$ as described in section 5.1. We have performed the measurements for different samples with varying thicknesses for the top and bottom Pt layer. It seems that varying the thickness of the bottom Pt layer changes the damping but we only observe this at low fields. At low fields the noise is too large to draw any conclusions. Varying the thickness of the top layer does not seem to influence the damping. We will now show measurements on a different substrate, which are of much better quality.

For reasons we explained in chapter 3, at some point we switched the substrate. We started to use substrates with a 100 nm oxide layer. It turned out that we obtained a signal with much less noise, which leads to more trustable fits of $\alpha$. We sputtered new samples on this substrate, again with varying Pt layer thicknesses and performed the same damping measurements as discussed before. The results are plotted in figure 5.5. In figure 5.5a we plotted the results for varying bottom layer and in figure 5.5b for varying top layer. We indeed observe that the signal has improved significantly.
5.4 Damping measurements with current induced spin torques

In chapter 2, macro-spin simulations revealed that a SHE induced spin current should be able to affect the effective $\alpha$ in Pt/CoB/Pt nano-wires, see figure 2.16. We argued that the largest effect should be present using an applied in-plane magnetic field in the $y$ direction. The main goal of this research was to investigate whether we could measure this effect using the TR-MOKE set-up to determine $\alpha$. In this final section we will discuss those current influenced damping measurements. During those measurements we used a continuous current. Therefore the maximum current density we can apply is limited by Joule heating. This means that we expect only small changes in $\alpha$. Therefore we needed to optimize the signal as much as possible. In this chapter we already discussed the influence of the pump-laser power and the substrate on the signal to noise ratio. In section 3.2.3 of chapter 3 we also discussed some of the other techniques used to improve the signal as much as possible. In this section we show the final measurements, obtained after all optimizations. We performed measurements on both a Pt(5)/CoB(0.7)/Pt(2) as well as a Pt(5)/CoB(0.7)/AlOx(0.89) nano-wire. During all those measurements an $H_{yz}$ (see figure 4.1b) field is used to maximize the

![Graph showing damping parameter $\alpha$ as a function of applied field for different Pt thicknesses on Pt/CoB/Pt films. During those measurements, $\Theta=82.5^\circ$. Sample fabricated on Si with a 10 nm SiOx top layer.](image)

We observe no dependency of $\alpha$ on the top Pt layer thickness. However, we observe a small dependency of $\alpha$ on the bottom Pt layer thickness. Because we do not observe an effect of the top layer thickness, we can conclude that the Pt layer thickness does not influence the damping. The reason a small dependency of the bottom layer is observed, is probably due to a difference in growth for the first layers.
CHAPTER 5. CURRENT INFLUENCED MAGNETIC DAMPING

Figure 5.5 – Damping as a function of the applied magnetic field. During those measurements, $\Theta = 82.5^\circ$. The error bar indicates the standard deviation in the fitting procedure. a) Measurements performed for varying bottom layer thickness b) Measurements performed for varying top layer thickness. $\Theta = 82.5^\circ$ during those measurements. Sample fabricated on Si with 100 nm SiOx top layer.

Before we performed measurements with an applied current, we measured $\alpha$ as a function of applied magnetic field on a Pt(5)/CoB(0.7)/Pt(2) nano-wire, like we did in the previous section for thin films, to characterize our system. We want to perform current assisted measurements with an applied field that is large enough to obtain the intrinsic value of $\alpha$, because in that case we obtain a homogeneous precession in the sample. The results are given in figure 5.6b. We decided to do the measurements using a field of 332 mT because at that field we have almost reached the intrinsic value of $\alpha$. If we use larger applied fields the signal quality started to decrease.

Subsequently we performed damping measurements while sending an electrical current through the nano-wire. In figure 5.6a we show the results of damping measurements performed on a 50 $\mu$m wide Pt(5)/CoB(0.7)/Pt(2) nano-wire with an applied current and an applied magnetic field of 332 mT under an angle of 80 degrees with the z axis. We repeated the measurement 10 times for each value of the current and plotted the average value of $\alpha$ over those 10 measurements. We also calculated the standard deviation and plotted it along with $\alpha$. We only observe random oscillations of $\alpha$, but no well defined trend as a function of the applied current. The maximum current density we plotted is $1.6 \cdot 10^{10} \text{A/m}^2$. From the macro-spin simulations explained in chapter 2, we expect a change in $\alpha$ of only 0.004 for this current density. However, the standard deviation in $\alpha$ in our measurements is around 0.005. Unfortunately, when we increase the current any further we observe non reproducible effects.
5.4. DAMPING MEASUREMENTS WITH CURRENT INDUCED SPIN TORQUES

Figure 5.6 – a) Gilbert damping parameter measured in a 50 µm Pt(5)/CoB(0.7)/Pt(2) nano-wire as a function of the applied current. For every current the measurement is repeated 10 times. We plotted the average value of $\alpha$ and the standard deviation. No effect from the current can be observed. A field strength of 332 mT is used. b) The damping parameter as a function of applied field. We observe the same behavior as for the thin films. During those measurements, $\Theta=80^\circ$.

in our measurements due to Joule heating. This means that with the obtained signal quality and the maximal possible current for which we can perform reproducible measurement, it is not possible to measure the SHE induced effect on the damping parameter in Pt/CoB/Pt nano-wires. We should reduce the standard deviation to be able to observe the effect.

We also performed measurements on a 50 µm wide Pt(5)/CoB(0.8)/AlOx(0.89) wire, using an applied magnetic field of 580 mT and an angle of 80 degrees with the z axis. The results are shown in figure 5.7a. Because this system has a large intrinsic damping parameter, we observe only one complete oscillation, see figure 5.8b. This makes the determination of $\alpha$ even less accurate as for the Pt/CoB/Pt system. When we take a look at figure 5.7a we observe a increase in $\alpha$ for small currents and an decreasing $\alpha$ for large currents, that is more or less symmetric with respect to the current direction. Although those damping measurements were less accurate, we clearly see an current induced effect when we plot the time evolution of the z component of the magnetization, see figure 5.8b. We plotted the trajectory for different values of the applied current density and normalized the different plots on the first peak, to be able to compare them. For high currents we observe that the second oscillation is better visible, which means that the damping should be smaller. Those measurements cannot be explained by the SHE or the Rashba field, so more research necessary.

Because of the high damping parameters in the studied systems, we were not able
Figure 5.7 – a) Damping as a function of current density measured in a 50 µm Pt(5)/CoB(0.8)/AlOx(0.89) nano-wire. The damping seems to reduce for higher currents, but no asymmetry between positive and negative current can be observed. During those measurements, $\Theta=80$° and the applied field is 580 mT. b) Damping parameter as a function of applied field for the Pt/CoB/AlOx system. During those measurements, $\Theta=80$°. Also the standard deviation is depicted.

Figure 5.8 – Magnetic signal from the second lock-in amplifier. The data is normalized on the first peak. During those measurements, $\Theta=80$° and the applied field was 580 mT.
to detect a current induced effect on the magnetic damping. Our measurements are mainly limited by Joule heating and the high damping parameter. So future research should focus on those two aspects. Some suggestions to improve the technique are discussed in chapter 6.
Chapter 6

Conclusions and outlook

This final chapter discusses the most important results obtained in this thesis and the conclusions we can draw from them. In the second section an outlook on possible future research is given. We will also take a look at an application of the SHE in magnetic memory devices.

6.1 Conclusions

Over the last years, the interest in current induced magnetization manipulation has increased continuously, due to its promising properties for future memory applications. In this work we investigated the influence of an electric current on the magnetization dynamics, in magnetic nano-wires. In particular, we choose to study the current-induced torques in perpendicular magnetic anisotropy materials (e.g., Pt/Co/Pt and Pt/Co/AlOx). Indeed, in these systems the presence of high spin-orbit coupling makes the current induced dynamics much more complex, as it was observed in the literature [16] and in our group. Two spin-orbit induced effects are proposed which could explain the observations: the spin Hall effect (SHE) and Rashba spin-orbit coupling. In this thesis we focused on characterizing the SHE in Pt/CoB/Pt systems, for which the Rashba contribution is neglectable. We used two types of measurements to investigate the SHE in those systems.

First we characterized the SHE using current assisted field induced switching measurements. From symmetry considerations and macro-spin simulations we expect the SHE to influence magnetic switching, when it is combined with an applied magnetic field in the x direction. This effect was indeed demonstrated by switching experiments performed on a TR-MOKE set-up. As predicted by the simulations, the field induced magnetic switching is assisted or opposed depending on both the current and the in-plane magnetic field direction. We also showed that we can switch the magnetization using only the spin Hall effect in combination with an applied in-plane magnetic field. Also in those measurements, a current induced effect with the symmetry of the
SHE is observed. However, we were not able to systematically study the SHE using those measurements, because the in-plane magnetic field brings the system in a poorly defined multi-domain state.

Subsequently, we investigated a new technique to study the SHE. In the correct configuration, the torque induced by the SHE can have a component along the direction of the damping torque. This means that we should be able to study SHE by measuring its effect on the $\alpha$. An analysis of the effect of the injected spin indicated that the effect is maximized when an in-plane field perpendicular to the wire is used. We performed macro-spin simulations which confirmed this current induced effect on $\alpha$. The simulations provided an estimation of the size of the current induced effect on $\alpha$.

The TR-MOKE set-up was used to investigate the current influenced magnetic damping. Those measurements are very complex for a couple of reasons. Out-of-plane magnetized Pt/CoB/Pt systems have a large damping, which make the determination of $\alpha$ less accurate. On top of that, Joule heating limits the current we are able to send through the nano-wires. This means that a high signal to noise ratio is desired to be able to detect those small changes in damping. Therefore, a lot of effort was put into the optimization of the set-up. With the obtained signal to noise ratio, we were not able to detect a current induced effect on $\alpha$. Further optimization is needed to use this technique to study the SHE in perpendicular magnetized systems.

### 6.2 Outlook and applications

This final section will give an outlook to future research and applications. As discussed in the previous section, further optimizations are needed to be able to study the effect of current induced spin-orbit torques on the magnetic damping. In the first part of this section, some suggestions their feasibility are discussed. We will conclude with an application of the SHE in magnetic memory.

The first limitation in our measurements was the maximum applied current we are able to send through the nano-wire. We used continuous currents, which leads to large Joule heating. We could increase the maximum current while keeping the Joule heating constant by using pulsed currents. Those current pulses have to be synchronized with the pump-beam, and last for at least the duration of the damped oscillation, which is typically 500 ps. Those current pulses could be generated by a pulse generator that can be synchronized with the laser, that is available in this group.

Another opportunity is to replace the CoB layer by a perpendicular magnetized material with a lower damping parameter. This would lead to a more accurate determination of $\alpha$. At this moment, an interesting system is investigated in this group: Pt/(Co/Ni)$_n$/Pt systems, were the Co/Ni combination is repeated $n$ times. When the correct Co an Ni thickness is used, this system has out-of-plane anisotropy for values of $n$ even above 6. Besides, this system has a damping parameter around 0.035, which is
very small compared to the Pt/CoB/Pt system [49]. Therefore it would be interesting to study current influenced damping in those systems. We performed a small study to make a rough estimate of how accurate we can determine $\alpha$ in this system. The standard deviation in the damping measurements performed on Pt/CoB/Pt is around $5 \cdot 10^{-3}$. Using this value, we could use macro-spin simulations, in which we introduced noise, in order to make a rough estimate of the expected standard deviation in damping measurements on Co/Ni systems. We found an standard deviation of $2.3 \cdot 10^{-4}$, 22 times smaller than for the measurement performed on Pt/CoB/Pt samples. Therefore, although this analysis is very rough and neglects a number of other important aspects, we think that this is a promising system to consider.

To conclude we take a look at a practical application of the SHE in a magnetic memory device. An important application of SHE induced STT could lie in the development of a new kind of MRAM. A MRAM element consists of two magnetic layers separated by a tunnel barrier. One of the magnetic layers is pinned, the second is free. In present MRAM devices, the free layer is switched by sending a current through the stack. The spin polarized current from the fixed layer entering the free layer is able to switch the free layer parallel or anti-parallel to the fixed layer, depending on the direction of the current. The relative direction of the free layer compared to the fixed layer determines the bit state, which is measured using the GMR effect. This switching process requires high currents flowing through the tunnel barrier. When we use an adjacent Pt layer to inject a spin current into the free layer (using the SHE), we could switch the magnetization without a current through the tunnel barrier. In this way, there is no need to send high currents through the fragile tunnel barrier, which therefore can be optimized for the measurement of the GMR.
Appendix A: Joule heating in nanowires: measurements and simulations

In this appendix we will take a closer look at the temperature simulations and measurements discussed in 4.3.1. In the first section we discuss the simulations. We performed finite element simulations using COMSOL multiphysics to study Joule heating in magnetic nano-wires. In the second section we discuss measurements performed to determine the temperature rise due to Joule heating.

Simulations

We used COMSOL multiphysics to simulate Joule heating in magnetic nano-wires. We gave the results in section 4.3.1. Here we will provide some extra details regarding those simulations. We assumed that in equilibrium, the temperature in the whole nano-wire is homogeneous and that all heat is dissipated through the substrate. This gives us the opportunity to model only the substrate. This is necessary, because the difference in dimensions of the substrate and the thickness of the nano-wire gives problems during the meshing process in COMSOL. We simulated a Si substrate with a 100nm AlOx layer on top of it. The performed simulations are 2d, which means that we assume an infinite long wire. We simulated a 2mm thick and 10 mm wide substrate. We keep the boundaries of the substrate at room temperature. A typical temperature distribution is given in figure 6.1. We determined the temperature of the wire a function of the applied current in the wire.

Measurements

We also performed measurements to determine the temperature of the nano-wires as a function of the current trough the wire. We performed the measurements on a Pt(5)/CoB(0.8)/AlOx(0.89) nano-wire The measurements make use of the tempera-
Figure 6.1 – Simulations of the temperature distribution in the substrate due to Joule heating in a nano-wire on top of the substrate. Simulated is a 2d Si substrate with a width of 10 mm and a thickness of 2 mm and a 100 nm SiOx layer on top. This layer is not visible in the figure because its very small in comparison with the size of the substrate.

We measured the resistance as a function of time for different values of the applied voltage. A typical measurement is shown in figure ???. The resistance will increase over time due to the heating of the wire. We see that it takes typically a few minutes for the system to reach its equilibrium temperature.

If we measure the change in resistance for a series of different applied voltages, we can calculate the temperature change as a function of applied voltage with the help of a relation between resistivity and temperature for Pt. In this calculation we assumed that the complete wire was made of Pt. The formula for the resistivity of Pt as a function of the temperature is taken from the COMSOL library [14]. The results are given in figure 4.14a.
6.2. OUTLOOK AND APPLICATIONS

Figure 6.2 – *Typical measurement of the resistance of the nano-wire as a function of time. It takes typically a few minutes for the wire to reach its equilibrium temperature.*
CHAPTER 6. CONCLUSIONS AND OUTLOOK
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