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Ultrafast laser-induced spin-transfer torque in a non-collinear magnetic bilayer

de Wit, R.R.J.C.

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Ultrafast laser-induced spin-transfer torque in a non-collinear magnetic bilayer

Ruud de Wit

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Supervisors:
MSc A.J. Schellekens
prof. dr. B. Koopmans
Abstract

Many proposed spintronics devices for writing data are based on the use of spin currents. These currents can switch magnetic bits by transfer of angular momentum from one magnetic layer to the other, exerting a torque on the magnetization. To switch magnetic devices efficiently with current-induced spin-transfer torque, new methods are being investigated to create intense and short spin current pulses. In this thesis, we report on the occurrence of spin transfer torque on unprecedented timescales by excitation of a ferromagnetic bilayer by femtosecond laser pulses.

When exciting a ferromagnetic layer with a femtosecond laser pulse, an ultrafast demagnetization is induced in this layer within a picosecond. This demagnetization has shown to be accompanied by the generation of a spin-polarized current on the same extremely short timescale. In this research, it is demonstrated that laser-induced spin currents are present in a CoNi/Cu/Co non-collinear magnetic bilayer and that these currents exert a spin-transfer torque on one of the magnetic layers, changing the orientation of its magnetization. To this end, we performed time-resolved laser-induced remagnetization measurements using TR-MOKE, from which laser-induced precessions of the magnetization in both magnetic layers are isolated. The properties of a precession, such as amplitude and phase, depend on the excitation mechanism of the precession. We developed macro-spin models for precessions induced by different mechanisms in both layers and compared the obtained precessional properties to these models. It is demonstrated that the in-plane magnetized layer of the bilayer is excited by a short majority spin current, coming from the out-of-plane magnetized layer. Furthermore, the amplitudes of the current-induced precessions are compared for Cu and Pt spacer layers. The induced precessions are found to be much smaller for Pt, as expected from the small spin diffusion length in Pt. Furthermore, we found that the spin transfer length in Pt is almost ten times smaller than in Cu.

Concluding, we have demonstrated for the first time that laser-induced spin currents can be used to change the orientation of the magnetization of one of the magnetic layers in a non-collinear bilayer.
## Contents

Abstract

1 Introduction 1
  1.1 Giant magnetoresistance ................................. 1
  1.2 Spin-transfer torque .................................... 3
  1.3 Laser-induced demagnetization .......................... 5
  1.4 This research ........................................... 7
  1.5 This thesis ............................................. 8

2 Theory 9
  2.1 Magnetism ............................................... 9
    2.1.1 Ferromagnetism ...................................... 9
    2.1.2 Anisotropy ......................................... 11
    2.1.3 Hysteresis .......................................... 14
  2.2 Magnetization dynamics ................................. 15
    2.2.1 The Landau-Lifshitz-Gilbert equation ............... 15
    2.2.2 Precessions: macro-spin calculations ............... 17
    2.2.3 Example: phase difference for short and long spin-
      transfer torques ....................................... 19
    2.2.4 Phenomenological description of remagnetization after
      laser excitation ....................................... 21
  2.3 Mechanisms for inducing a spin current .................. 24
    2.3.1 Superdiffusive transport ........................... 24
    2.3.2 Spin-dependent Seebeck effect ....................... 26

3 Experimental 29
  3.1 Phenomenological description of MOKE ................... 29
  3.2 Time-Resolved Magneto-Optical Kerr Effect ............... 31
    3.2.1 Simple TR-MOKE setup ................................ 31
    3.2.2 Improved TR-MOKE setup ............................. 34
CONTENTS

4 Results 37
  4.1 Static magnetic properties of the bilayer 37
  4.2 Acquisition of precessional properties 41
  4.3 Distinguishing the magnetic layers 42
  4.4 Excitation mechanisms of laser-induced precessions 44
    4.4.1 Phase measurement 45
    4.4.2 Comparing phases: possible excitation mechanisms 46
  4.5 Field-dependent amplitude 47
  4.6 Angle-dependent measurements 49
  4.7 Comparing precessions in copper and platinum spacer layers 51
  4.8 Quantification of spin currents 53
  4.9 Microscopic origins of the spin currents 56

5 Conclusion and outlook 59
  5.1 Conclusions 59
  5.2 Outlook 60

A The out-of-plane layer 61
  A.1 Equilibrium state before laser excitation 61
  A.2 Field-dependent frequency 63
  A.3 Field-dependent amplitude of an anisotropy pulse-induced precession 66

B The in-plane layer 69
  B.1 Field-dependent frequency 69
  B.2 Field-dependent amplitude of a spin current-induced precession 71
    B.2.1 Derivation 71
    B.2.2 Approximation for small spin currents 73

C Influence of the applied field angle 77

D Quantification of the spin currents 79

Bibliography 81
Chapter 1

Introduction

In this chapter, the scientific background and the main goal of the work presented in this report are treated. We will start with a brief introduction to the research field of spintronics. After that, the main experiment is presented.

1.1 Giant magnetoresistance

Although the concept of spin was already suggested by Uhlenbeck and Goudsmit in 1926 [1], we had to wait another 70 years for the field of spintronics to be started. It started all with Albert Fert [2] and Peter Grünberg [3] in respectively 1988 and 1989. They experimentally demonstrated that an electrical current in a metal can be divided into two separate spin channels (up and down). This was already suggested in 1936 by Sir Neville Mott [4]. Fert and Grünberg independently discovered that the electrical resistance of these spin channels is not necessarily equal. In particular, the resistance of a ferromagnet turned out be completely different for majority spins and minority spins. Fert and Grünberg demonstrated this by the effect called giant magnetoresistance (GMR). It became the cornerstone of spintronics and Fert and Grünberg were awarded a Nobel prize for their discovery in 2007.

GMR is based on the fact, that the total electrical resistance of two ferromagnetic layers depends on how the magnetization in each layer is oriented. This is shown in Figure 1.1. An electrical current $I$ flows through two ferromagnetic layers with a parallel or anti-parallel oriented magnetization, depicted by the black arrows. The current is regarded as two spin channels, experiencing a different resistance. For example, the up-spins experience a high resistance $R_H$ in the upward magnetized layer and a low resistance $R_L$ in the downward magnetized layer. Using the equivalent circuits on the right in Figure 1.1, the total resistance can be calculated for both orientations.
CHAPTER 1. INTRODUCTION

The resistance for the parallel (P) and the anti-parallel (AP) orientation are respectively given by:

\[
\frac{1}{R_P} = \frac{1}{R_L} + \frac{1}{R_H} = \frac{2}{R_L + R_H}, \quad (1.1)
\]

\[
\frac{1}{R_{AP}} = \frac{1}{R_L} + \frac{1}{R_H} = \frac{R_L + R_H}{2R_L R_H}. \quad (1.2)
\]

Thus, the resistance depends on the orientation of the magnetization of magnetic layers. The relative difference between both resistances is defined by the magnetoresistance

\[
\frac{\Delta R}{R} = \frac{R_{AP} - R_P}{R_P} = \frac{(R_H - R_L)^2}{4R_L R_H}. \quad (1.3)
\]

Soon after the introduction by Fert and Grünberg, huge magnetoresistances of tens of percents were observed and the effect was immediately labelled as giant magnetoresistance, or GMR. The GMR effect opened new possibilities for developing sensitive magnetic field sensors. In 1997, IBM was the first to introduce a read head of a hard drive, based on GMR. By simply measuring the magnetoresistance of the GMR sensor, the direction of the magnetic field in the bit - representing 0 and 1 - could be deduced.

\[\text{Figure 1.1: Giant magnetresistance. The total electrical resistance of two ferromagnetic layers depends on the orientation of the magnetization of the layers. Courtesy to Swagten [5].}\]
1.2 Spin-transfer torque

The read head of IBM introduced a new way of reading data using GMR. However, the physics governing GMR also opened up new possibilities for writing data. In 1996, Slonczewski [6] predicted a new way of exciting the magnetic state of a ferromagnet by introducing a phenomenon called current-driven spin transfer.

Consider an electrical current flowing through the multilayer of Figure 1.2. The blue layers are ferromagnets and the yellow layers are nonmagnetic spacer layers. The bottom is a pinned ferromagnet: a magnetic layer of which the magnetization is fixed. The magnetization of the top magnetic layer can freely change direction. Slonczewski predicted that the bottom magnetic polarizes the current in the direction of $\vec{S}_1$. Consequently, the current is polarized when it enters the spacer layer. This phenomenon is called spin injection. If the thickness of the spacer layer is less than the spin-diffusion length (about 100 nm), the current will still be spin-polarized as it reaches the second ferromagnet. In this layer, spin transfer takes place between the spin-polarized current and the spins in the magnetic layer. Consequently, the magnetization, initially pointing along $\vec{S}_2$, is tilted towards $\vec{S}_1$ by the so-called spin transfer torque.

![Figure 1.2: Spin-transfer torque in a multilayer. The blue layers are ferromagnets and the yellow layers are nonmagnetic (NM) spacer layers. A spin-polarized current, coming from the bottom layer, exerts a spin-transfer torque on the magnetization of the top magnetic layer.](image-url)
This means that the magnetic state of a ferromagnet can be changed by a spin-polarized current. Slonczewski already predicted that spin-transfer torque can, in potential, be used for the switching of the magnetization of ferromagnets. Hence, current-driven magnetization switching is a promising new way of writing data. Probably the most revolutionary concept based on storing data with currents is the racetrack memory of Parkin et al. [7]. In this device, data is stored as magnetic domains, which are moved through the device by currents.

Another recent example of current-driven magnetization switching is the research of Mangin et al. [8]. In this paper, current-driven magnetization switching is demonstrated in a so-called spin valve, depicted in Figure 1.3. The Co/Ni layers (in green) are the ferromagnetic layers. In this stack, both magnetic layers are magnetized perpendicularly to the film plane. An electrical current will become perpendicularly spin-polarized by the bottom ferromagnetic layer. Subsequently, this current exerts a spin-transfer torque on the magnetization of the top magnet. If this torque is large and long enough, the magnetization is switched.

![Figure 1.3: Magnetization switching in a perpendicular spin valve. The green layers are ferromagnetic layers. Both ferromagnetic layers are magnetized perpendicular to the film plane. The magnetization of the top magnet is switched by a perpendicularly spin-polarized current, coming from the bottom layer. Courtesy to Mangin et al. [8].](image)

Concluding, the magnetization switching can be induced by spin-transfer torque, caused by spin-polarized currents. In the example above, these spin-polarized currents are achieved by sending an electrical current through the spin valve. However, in the research presented in this thesis, spin-polarized currents are created on extremely short timescales by femtosecond laser pulse excitation. This will be discussed in the next section.
1.3 Laser-induced demagnetization

When a ferromagnet is excited by a femtosecond laser pulse, its magnetization drops rapidly. It was demonstrated in 1996 by a pioneering experiment that such a demagnetization takes place at subpicosecond timescales [9]. The original measurement is shown in Figure 1.4. This fast demagnetization was assigned to ultrafast heating of the spin system in the ferromagnet, lowering the magnetization.

![Figure 1.4: Pioneering measurement of laser-induced demagnetization. The magnetization decreases within a the picosecond upon excitation by a femtosecond laser pulse. Courtesy to Beauerpiaire et al. [9]](image)

Other work in the literature has indeed demonstrated laser-induced demagnetization in a ferromagnet within 300 fs [10, 11]. Furthermore, it is experimentally demonstrated that the magnetization of multilayer ferromagnets can be reversed in a reproducible manner by a femtosecond laser pulse [12]. Experiments on Co/Pt multilayers demonstrate that interlayer transfer of spin angular momentum is accompanied by an increase of the total demagnetization by almost 25 % [13].

The scientific community is still debating about the origin of this ultrafast demagnetization. It is unclear how laser excitation can couple so quickly to the spins. A model based on superdiffusive spin transport is proposed by Battiato et al. [14]. Recent experiments by Rudolf et al. [15] seem to confirm this model and thus the presence of laser-induced spin currents. We will treat the latter research briefly for illustrative purposes.
CHAPTER 1. INTRODUCTION

In this experiment, a Ni/Ru/Fe trilayer is excited with a laser pulse. In this sample, the Ni and Fe layers can be ferro- or antiferromagnetically coupled. It is demonstrated that laser-induced demagnetization of the Ni layer enhances the magnetization of the Fe layer when the magnetizations of both layers are initially aligned parallel. Simulations of the experiments are shown in Figure 1.5. Rudolf et al. [15] explain this by a laser-induced spin current between the layers.

So concluding, in our search for faster data storage devices, pulsed laser excitation of ferromagnets seems interesting, as it provides access to the ultimate timescale for magnetization dynamics. Furthermore, the fact that this ultrafast demagnetization is accompanied by the generation of spin-polarized currents opens up the possibility of current-induced switching of magnetic bits. It is believed that the excited electrons in the ferromagnetic layer travel in random directions through the material and, in particular, out of the material [14, 15]. The mean free path of these electrons is larger for majority spins than for minority spins. A current flowing out of the magnet will thus be spin-polarized, which is shown in Figure 1.6. Furthermore, this spin-polarized excited electron diffusion takes place in the superdiffusive regime, meaning that spin transport occurs semi-ballistically.

Concluding, a superdiffusive spin-polarized current can be generated on subpicosecond timescales by femtosecond pulsed laser excitation. In the next section we discuss the main goal of this report, namely the realization of spin-transfer torque on the ultimate timescale by generating laser-induced superdiffusive spin currents.
1.4 This research

In this thesis, we attempt to change the orientation of a magnetic layer by femtosecond laser-induced superdiffusive spin currents. The main concept of our research is schematically depicted in Figure 1.7.

The cartoon in Figure 1.7 illustrates how ultrafast laser-induced spin currents generate precessions in a non-collinear magnetic bilayer. This is explained by a sample consisting of an out-of-plane magnetized bottom layer and an in-plane magnetized top layer, divided by a conducting spacer layer. We hit this sample with an extremely short (\( \sim \) fs) laser pulse at \( t = 0 \). This heats the sample, creating high-energetic electrons. Subsequently, a spin-polarized current will possibly be generated. If so, both layers experience a spin-polarized current coming from the other layer. These currents are depicted by the black arrows in Figure 1.7. Since these layers are non-collinear, these currents exert a spin-transfer torque on the magnetization of the other layer, tilting the magnetization out of its easy axis. The magnetization then relaxes back to the equilibrium state by means of a damped precession, which is an elementary spin dynamics process. It will be described in section 2.2.1.

The goal of this research is to prove the existence of the ultrafast laser-induced spin currents discussed in the above. To this end, precessions of the magnetization are measured. These precessions could be the consequence of a spin-transfer torque, exerted by a spin current. However, precessions of the magnetization can also be induced by other excitation mechanisms. These mechanisms can be distinguished by looking at the properties of the induced precessions, such as the amplitude and phase and comparing the results to a macro-spin model.
CHAPTER 1. INTRODUCTION

1.5 This thesis

This thesis contains the following sections:

- Chapter 2 - Theory: important properties of ferromagnetic materials are treated and laser-induced magnetization dynamics in such materials are presented. Furthermore, two excitation mechanisms for inducing a spin-polarized current are discussed.

- Chapter 3 - Experimental: the TR-MOKE setup, with which the experiments are performed, is explained. Furthermore, the physics of MOKE are discussed.

- Chapter 4 - Results: the results are presented and compared to macrospin models. The presence of a laser-induced spin-transfer torque-induced precession of the magnetization in a non-collinear magnetic bilayer is demonstrated.

- Chapter 5 - Conclusions and discussion: conclusions are drawn from the measurements and the obtained results are discussed.
Chapter 2

Theory

In this thesis, laser-induced precessions of the magnetization of ferromagnetic layers are measured and compared to macro-spin models. We aim on demonstrating that these precessions are induced by spin-transfer torque, caused by a laser-induced spin current. In this chapter, we introduce the used theoretical concepts. First, we will give a brief introduction to ferromagnetism. Next, the dynamics of a laser-induced precession of the magnetization are treated. Finally, we will distinguish two excitation mechanism for inducing a spin-polarized current, which both can be present during ultrafast laser-induced demagnetization.

2.1 Magnetism

The measurements in this thesis are performed on magnetic thin films. Here, we introduce the concept of spin to explain why some materials are permanently magnetized. This phenomenon is called ferromagnetism. Furthermore, we explain some important properties of these ferromagnetic materials which are crucial for the interpretation of the experiments presented in this thesis, such as magnetic anisotropy and magnetic hysteresis.

2.1.1 Ferromagnetism

Magnetism originates from physics at the atomic level. It is based on the fact, that every electron has its own magnetic moment. We consider a single electron orbiting around a nucleus with a speed $v$ and at a distance $r$, see Figure 2.1. This electron corresponds to a current $\mathbf{I}$, flowing in the opposite direction. Now, it is well-known from Maxwell electrodynamics that a rotating current induces a magnetic moment which is proportional to $-I \mathbf{A} \times \mathbf{z}$. It
can be shown that this magnetic moment is related to the angular momentum of the electron \( \vec{L} = m_e v r \hat{e}_z \) by [5]

\[
\vec{\mu}_L = \gamma_L \vec{L},
\]  

(2.1)

in which \( \gamma_L \) is called the gyromagnetic ratio. Hence, an orbiting electron possesses an extrinsic magnetic moment because it carries charge.

![Figure 2.1: An electron orbiting around a nucleus. The corresponding current induces a magnetic moment \( \mu \). Courtesy to Swagten [5].](image)

Besides the magnetic moment due to the charge, each electron carries another property which induces a contribution to the magnetic moment: spin. This quantumphysical phenomenon can be seen as the self-rotation of an electron around its own axis, although the electron is in fact a structureless point particle. The concept of spin was first suggested by Uhlenbeck and Goudsmit in 1926 [1].

In addition to the angular momentum \( \vec{L} \) of the orbiting electron, each electron also carries intrinsic angular momentum \( \vec{S} \) due to spin. The most striking property of spin is that \( \vec{S} \) comes in discrete values. In the case of an electron, there are only two possible states: spin up and spin down. The magnetic moment corresponding to spin can be written as [17, 18]

\[
\vec{\mu}_S = \gamma_S \vec{S},
\]  

(2.2)

in which \( \gamma_S \) is the gyromagnetic ratio for spin [5]. The total magnetic moment of an electron orbiting around a nucleus is now simply given by

\[
\vec{\mu} = \vec{\mu}_L + \vec{\mu}_S.
\]  

(2.3)
Thus, every electron has its own magnetic moment \( \vec{\mu} \). The macroscopic magnetization of a material is now defined as the vectorial sum of all magnetic moments per unit of volume:

\[
\vec{M} = \Sigma \vec{\mu} / V,
\]

in which \( V \) is the total volume of the object. In most materials, the direction of the magnetic moments is random. Hence, the average magnetic moment is zero and these materials are intrinsically non-magnetic. An example is the class of paramagnet materials, shown on the left in Figure 2.2 for a finite temperature. The arrows depict the random direction of the magnetic moments. Paramagnetic materials can only be magnetized by an externally applied magnetic field. As soon as this field is removed, these materials become non-magnetic again.

However, some materials contain a net magnetic moment \( \vec{M} \), even without the presence of a magnetic field. These materials are called ferromagnets. Well-known ferromagnetic materials are Fe, Ni and Co. The structure of the magnetic moments in a ferromagnet is depicted on the right side of Figure 2.2. In a ferromagnetic material, it is energetically the most favorable for the spins to align in one direction. This collective ordering is due to exchange interaction between the spins on the atomic level. However, we will not go into further detail. For further reading, see Swagten [5].

\[\text{Figure 2.2: Orientation of magnetic moments in a paramagnetic material (left) and a ferromagnetic material (right). The paramagnet has no net magnetic moment. The magnetic moments in the ferromagnet are aligned and add up to an intrinsic magnetization. Courtesy to Swagten [5].}\]

### 2.1.2 Anisotropy

In this thesis, measurements are performed on ferromagnetic thin films. The direction of the magnetization in such films is not arbitrary: the magnetic
field in a magnetic thin film prefers to point along a so-called easy axis. The phenomenon causing this preferred directing is called magnetic anisotropy. Three types of anisotropy play a role in our research: shape anisotropy, crystalline anisotropy and surface anisotropy. We start with the explanation of shape anisotropy.

**Shape anisotropy**

We have seen in section 2.1.1 that the average direction of all spins in a material decides whether the material is intrinsically ferromagnetic or not. In a ferromagnetic material, all spins point in one direction (for zero temperature), causing a magnetization $\mathbf{M}$ in this direction. Now, let us consider a finite amount of ferromagnetically coupled spins, represented in Figure 2.3 by bar magnets. It can be seen that free magnetic poles occur at the open ends on both sides. These poles induce a so-called demagnetization field $\mathbf{H}_d$, which opposes $\mathbf{M}$. The demagnetization field microscopically originates from dipole coupling. The size of $\mathbf{H}_d$ depends on the distance between the poles. This means that, for an arbitrarily shaped piece of material, the demagnetization field is different in every direction. Therefore, the demagnetization is in general written as $[5]$

$$\mathbf{H}_d = -\mathbf{N} \mathbf{M},$$

(2.5)

with $\mathbf{N}$ the shape-dependent demagnetization tensor. When we apply the above to the magnetic thin films used in our research (Figure 2.4), we immediately recognize two limits. Along the long axis of the film (the x-axis) $\mathbf{H}_d = 0$ , i.e. $N_x = 0$, since the ends on both sides are far apart. In the perpendicular direction (the z-axis), however, $N_z = 1 [5]$. Consequently, the demagnetization field is $\mathbf{H}_d = -\mathbf{M}_z$ in the perpendicular direction. This
means that it is unfavorable for the magnetization to point out-of-plane. Hence, the thin magnetic film of Figure 2.4 is in-plane magnetized.

![Figure 2.4: Magnetic anisotropy in a ferromagnetic thin film. A cross section of the thin film is shown. a) There is no demagnetization field when the magnetization points in the film plane, since both ends are too far apart. b) In the perpendicular direction, the demagnetization field is equally large and opposite to the magnetization. Therefore, the magnetization of the film has no out-of-plane component in the ground state.](image)

### Surface and crystalline anisotropy

If it were for the shape anisotropy, all magnetic thin films would be in-plane magnetized in the ground state. We do, however, know that this is not the case: thin films can also be magnetized out-of-plane. This is possible due to the existence of surface and crystalline anisotropy. We start with the latter.

An orbiting spin in a crystal is, in principle, not influenced by the surrounding lattice. However, if we take a reference frame in which the spin is motionless, the surrounding charges of the crystal orbit around this spin. Hence, the spin experiences a magnetic field from the surrounding lattice. This is called spin-orbit coupling. Due to this coupling, the spin - and thus the magnetization - prefers to be directed along certain axes of the crystal [5]. This phenomenon is called crystalline anisotropy. The easy-axes can be different for every material. Magneto-crystalline anisotropy is described by a volume-dependent term $K_V$.

There is another contribution to the anisotropy, which is caused by the surface anisotropy. This surface anisotropy has the same origin as the crystalline anisotropy, however, symmetry is not broken by the crystal but by an interface. Surface anisotropy is described by the constant $K_S$.

Both types of anisotropy can be combined by defining an effective shape-dependent constant [19]

$$K_{\text{eff}} = K_V + \frac{2K_S}{d},$$

(2.6)
with \( d \) the thickness of the film. The factor 2 originates from the two surfaces of the plane, assuming that the considered film is sandwiched between identical layers. In thin magnetic films the anisotropy energy can be written as [20, 19]

\[
E_{\text{ani}} = -K_{\text{eff}} \sin^2 \theta, 
\]

in which \( \theta \) is the angle between the magnetization and the film plane.

Finally, the effective anisotropy field then reads

\[
H_{\text{ani}} = -\frac{1}{\mu_0} \nabla E_{\text{ani}} = \frac{2K_{\text{eff}}}{\mu_0 M_{\text{sat}}} \sin \theta = \frac{2K_{\text{eff}}}{\mu_0 M_{\text{sat}}} \frac{M_z}{M_{\text{sat}}} = H_K \frac{M_z}{M_{\text{sat}}},
\]

with \( M_{\text{sat}} \) the saturation magnetization of the magnetic thin film, \( M_z \) the perpendicular component of the magnetization, \( H_K = 2K_{\text{eff}}/\mu_0 M_{\text{sat}} \) and \( \nabla E_{\text{ani}} \) the gradient of the anisotropy energy.

So concluding, we have seen three types of magnetic anisotropy: shape, surface and magneto-crystalline anisotropy. Together, they define the magnetic easy-axis and thus the preferential direction of the magnetization in a ferromagnetic layer.

### 2.1.3 Hysteresis

In this research, measurements are performed on non-collinear magnetic bilayers. The magnetostatic behavior of the magnetic layers in these bilayers is important as it shows how the magnetization in both layers reacts upon the application of a magnetic field. This behavior is described by a phenomenon called hysteresis. We will explain it here.

We have seen in section 2.1.1 that a ferromagnetic material is intrinsically magnetic, even without the presence of an external magnetic field. The magnetization is pointing along the so-called easy-axis of the magnetic layer. However, in the case of, for example, an out-of-plane magnetized layer, the magnetization can either point upward or downward. The system is in an energy minimum in both cases. To switch the magnetization, an external magnetic field has to be applied in the opposite direction. After switching off the field, the magnetization remains in its new state. This phenomenon is called hysteresis.

Figure 2.5a shows the perpendicular component of the magnetization of an out-of-plane magnetized ferromagnetic layer as a function of a vertically applied field. Such a plot is called a hysteresis loop. The coercivity \( H_c \) is the applied magnetic field strength at which the magnetization (the red arrow) switches. The black arrows show the field sweep direction. Figure 2.5b shows the same graph for an in-plane magnetized layer. In this case,
the magnetization is only tilted out of the film plane without switching, since the applied field is directed perpendicular to the magnetic easy-axis.

To conclude, the presented hysteresis loops show the magnetostatic behavior of layers, magnetized in different directions. We will compare these loops to our measurements to investigate the magnetostatic behavior of our samples.

2.2 Magnetization dynamics

In this thesis, the properties of laser-induced precessions of the magnetization in ferromagnetic layers are investigated. These properties are deduced from the magnetization dynamics after pulsed laser excitation. The dynamics of the magnetization of an excited ferromagnetic layer are treated in this section. We start with the dynamics of a precession of the magnetization, described by the Landau-Lifshitz-Gilbert equation. Then, it is shown how the properties of these precessions can be deduced from macro-spin models. After that, we derive a phenomenologically expression for the remagnetization of a magnetic layer after excitation by a femtosecond laser pulse.

2.2.1 The Landau-Lifshitz-Gilbert equation

The precessional motion of the magnetization $\vec{M}$ of an ensemble of spins is an elementary spin dynamics process. The dynamics of a precession of $\vec{M}$ are described by the Landau-Lifshitz-Gilbert (LLG) equation [21],

$$\frac{d\vec{M}}{dt} = -\gamma \mu_0 (\vec{M} \times \vec{H}_{\text{eff}}) + \frac{\alpha}{M}(\vec{M} \times \frac{d\vec{M}}{dt}), \tag{2.9}$$
in which $\gamma = 1, 76 \cdot 10^{11} \text{ rad/(s-T)}$ is the gyromagnetic ratio and $\mu_0 = 4\pi \cdot 10^{-7} \text{ T-m/A}$ is the vacuum permeability. A precession occurs when an effective field $\mathbf{H}_{\text{eff}}$ is present which is not parallel to $\mathbf{M}$. This is described in the first term on the right in equation 2.9. The effective field can contain many contributions different from the applied field. The second term on the right in equation 2.9 represents the damping of the precession due to energy dissipation. This damping is characterized by the Gilbert constant $\alpha$. An example of a damped precession of the magnetization around an effective field is illustrated in Figure 2.6.

![Figure 2.6: A typical damped precession of the magnetization around an effective field, governed by the LLG equation.](image)

In the research presented in this thesis, it is expected that precessions will be induced by spin-transfer torque, exerted by spin-polarized currents. Therefore, an extra term must be introduced in equation 2.9. The LLG equation then reads [22]

$$\frac{d\mathbf{M}}{dt} = -\gamma \mu_0 (\mathbf{M} \times \mathbf{H}_{\text{eff}}) + \frac{\alpha}{M} \left( \mathbf{M} \times \frac{d\mathbf{M}}{dt} \right) + \tau_{\text{STT}} (\mathbf{M} \times \mathbf{\sigma} \times \mathbf{M}), \quad (2.10)$$

in which $\tau_{\text{STT}}$ is proportional to the spin current and $\mathbf{\sigma}$ is the polarization direction of the spin current. It can be seen from equation 2.10 that the total effective field around which the magnetization precesses is now given by

$$\mathbf{H}_{\text{eff,STT}} = \mathbf{H}_{\text{eff}} + \mathbf{M} \times \mathbf{\sigma}. \quad (2.11)$$
CHAPTER 2. THEORY

After excitation, the magnetization will subsequently relax back to equilibrium by a precessional motion. By studying the details of this precessional motion it is possible to determine the microscopic origin of the excitation mechanism. Unfortunately, multiple excitation mechanisms could be triggered due to a femtosecond laser pulse, troubling the interpretation of the experiments. But, since the dynamics of the induced precessions differ for every kind of excitation mechanism, they can be distinguished. In this thesis, the frequencies, amplitudes and phases of measured precessions are analyzed, as they form a signature of the mechanism that caused the precession.

In the next section we will derive analytical expressions for the precessional dynamics of the LLG equation, which can be used for comparison to the measurements presented in Chapter 4.

2.2.2 Precessions: macro-spin calculations

The main goal of the research presented in this thesis is the demonstration of ultrafast spin-transfer torque after excitation of a non-collinear magnetic bilayer. We do so by analyzing the laser-induced precession due to spin-transfer torque in such a bilayer. However, precessions of the magnetization can also be induced by other excitation mechanisms. We stated in section 1.4 that the different excitation mechanisms can be distinguished by looking at the properties of the induced precessions, such as the amplitude and phase. In this thesis, these properties are measured as a function of an applied magnetic field and compared to macro-spin models to deduce which mechanisms excited the magnetic layers of the bilayer. The derivations of expressions for the precessional frequencies and amplitudes are presented in this section.

The expressions are all derived from macro-spin models. The derivations of all the expressions can, in full detail, be found in the Appendices A and B. In this section, we present a brief overview of the derivations, and show the main results.

As an example of how the characteristics of the induced precessions can be derived from the LLG equation as a function of applied field, we show a truncated derivation of an expression for the field-dependent frequency in an in-plane magnetized layer. The derivation can be found in more detail in Appendix B.1. The used model is illustrated in Figure 2.7. We use an approach, similar to the master thesis of Kuiper [23]. The magnetization \( \vec{M}_{\text{sat}} \) is initially pinned in the positive y-direction by an in-plane applied field \( \vec{H}_{\text{app}} \). We now assume that \( \vec{M} \) is tilted out-of-plane over a small angle \( \theta \) by an arbitrary excitation mechanism. This induces a precession, governed by the LLG equation. The field-dependent frequency is derived by solving equation 2.9.
CHAPTER 2. THEORY

The effective field $\vec{H}_{\text{eff}}$ is the sum of the anisotropy field $\vec{H}_d = -M_z$ (see section 2.1.2) and the applied field $\vec{H}_{\text{app}}$. Since $\theta$ is small, the demagnetization field is much smaller than the external magnetic field and the effective field approximately points along the $y$-axis. Hence, $\vec{M}$ is assumed to precess in the $(x,z)$-plane. We assume the magnetization to be initially tilted vertically upward and choose the ansatz

$$\vec{M} = \begin{pmatrix} M_x \\ M_y \\ M_z \end{pmatrix} = \begin{pmatrix} A_x \sin(\omega t) \\ M_{\text{sat}} \\ A_z \cos(\omega t) \end{pmatrix},$$

(2.12)

in which $M_{\text{sat}}$ is the saturation magnetization, $A_{x,z}$ are amplitudes, $\omega$ is the frequency and $t$ is the time after the initial perturbation. The effective field is given by

$$\vec{H}_{\text{eff}} = \begin{pmatrix} 0 \\ H_{\text{app}} \\ -M_z \end{pmatrix}.$$  

(2.13)

The equations 2.12 and 2.13 are now substituted into the LLG equation 2.9. Damping is not included in the analysis, since the Gilbert constant $\alpha$ is rather small [23]. This yields a set of three equations, which can be solved to find the frequency $\omega$. Inserting both expressions into the LLG equation 2.9 yields for the $x$- and $z$-direction

$$A_x \omega = \gamma \mu_0 A_z [M_{\text{sat}} + H_{\text{app}}],$$

(2.14a)
CHAPTER 2. THEORY

\[ A_x \omega = \gamma \mu_0 H_{\text{app}} A_x. \]  

(2.14b)

Solving the system yields

\[ f_{ip} = \frac{\gamma}{2\pi} \sqrt{B_{\text{app}}(B_{\text{app}} + \mu_0 M_{\text{sat}})}, \]  

(2.15)

where \( B_{\text{app}} = \mu_0 H_{\text{app}} \) is the applied field in T and \( f_{ip} \) is the frequency in Hz. Equation 2.15 is the Kittel equation for the frequency of a precession of the magnetization in an in-plane magnetized layer [24].

A expression for the field-dependent frequency is also derived for an out-of-plane magnetized layer in the same way. We came across some inconsistencies in the derivation of this expression in Kuiper [23]. These were corrected. The expression is given by

\[ \omega = \begin{cases} 
\gamma \sqrt{(B_{\text{ani}}^2 - B_{\text{app}}^2)} & , B_{\text{app}} < B_{\text{ani}} \\
\gamma \sqrt{(B_{\text{app}} - B_{\text{ani}})B_{\text{app}}} & , B_{\text{app}} > B_{\text{ani}} 
\end{cases}, \]  

(2.16)

in which \( B_{\text{ani}} \) is the perpendicular anisotropy field in T.

Furthermore, we derived expressions for the amplitude of precessions, induced by different excitation mechanisms in both layers of the bilayer. These expression are for the in-plane layer and the out-of-plane layer given by respectively

\[ A_{z,ip}^{\text{STT}} \approx C_{ip} \left(1 - \frac{B_{\text{app}}}{\mu_0 M_{\text{sat}}} \right), \]  

(2.17)

\[ A_{z,oop}^{\text{K}} \approx C_{oop} B_{\text{app}}^2, \]  

(2.18)

where \( C_{ip} \) and \( C_{oop} \) are arbitrary constants. The superscripts STT and K indicate precessions induced by respectively a spin-transfer torque and an anisotropy pulse.

In chapter 4, we present measurements of frequencies and amplitudes for varying applied fields. These measurements are compared to the expressions presented in this section.

2.2.3 Example: phase difference for short and long spin-transfer torques

Precessions of the magnetization can be induced by many excitation mechanisms. An effective way of excluding a lot of these mechanisms before proceeding with the macro-spin analyses is by looking at the phases of the obtained precessions. In this section, we show an example of how the phases...
are obtained from a qualitative analysis of precessions induced by different mechanisms.

**Figure 2.8:** Magnetization dynamics of spin-transfer torque-induced precessions in an in-plane magnetized layer for short and long pulses. The torque is induced by a minority spin current in this example. A magnetic field is applied parallel to the magnetic easy-axis. a) Magnetization dynamics for the first picoseconds after laser excitation. The spin current tilts the effective field towards the x-axis and the magnetization starts precessing around it. The precession in the (x,z)-plane is shown on the right. The little black dot depicts the position the magnetization after a few picoseconds. b) Magnetization dynamics for longer timescales. For a long pulse (drawn in red), the magnetization keeps precessing around the unchanged effective field. For a short pulse (drawn in blue), the spin current vanishes and the magnetization starts precessing around the new effective field, pointing along the y-axis again. c) Schematic plot of the z-component of the magnetization as a function of the time after laser excitation in picoseconds. The phases are 0 for a long pulse and π/2 for a short pulse.
In this example, it is shown that the phase of induced precessions depends on the duration of the excitation mechanism. This duration is compared to the typical evolution time of the precessions, which is in the order of 100 ps (obtained from measurements). We define two regimes: a short excitation pulse lasts much shorter than 100 ps and a long pulse lasts much longer.

As an example, we compare the phases of spin-transfer torque-induced precessions in the in-plane top layer for short and long pulses. We consider a laser-induced spin-transfer torque, caused by a current of minority spins, coming from the bottom layer. This current is $\uparrow$-polarized for the used configuration.\[\downarrow\]

The movement of the magnetization is depicted in Figure 2.8a for the first picoseconds after laser excitation. The effective field $\vec{H}_{\text{eff}}$ and the magnetization $\vec{M}$ initially point along the applied field $\vec{H}_{\text{app}}$. The $\uparrow$-polarized spin current $\vec{\sigma}$ tilts $\vec{H}_{\text{eff}}$ in the direction of $\vec{M} \times \vec{\sigma}$ (the x-axis), according to equation 2.11. Therefore, $\vec{M}$ starts precessing around the new effective field. The illustration on the right shows this precession in the (x,z)-plane. The little black dot depicts the position of $\vec{M}$ after a few picoseconds.

Figure 2.8b treats the magnetization dynamics for longer timescales. For a long excitation pulse (drawn in red), $\vec{H}_{\text{eff}}$ is unchanged with respect to the situation of Figure 2.8a and, therefore, $\vec{M}$ keeps precessing around this field. However, for a short pulse (drawn in blue), the spin current vanishes and $\vec{H}_{\text{eff}}$ points along $\vec{H}_{\text{app}}$ again. In this case, $\vec{M}$ will precess around the y-axis after the first picoseconds.

In Figure 2.8c, the z-component of the magnetization $M_z$ is schematically plotted as a function of the time after laser excitation in picoseconds. This component is also measured in our experiments. The phases, defined by equation 4.1, are 0 for a long pulse and $\pi/2$ for a short pulse. So, it can be concluded that the duration of the excitation pulse influences the phase of an induced precession.

### 2.2.4 Phenomenological description of remagnetization after laser excitation

We have shown in the previous sections that properties of precessions of the magnetization in magnetic layers are analyzed in this research. To this end, time-resolved remagnetization measurements are performed on magnetic layers after excitation by a femtosecond laser pulse. Subsequently, the precessional properties are deduced from these measurements by fitting them to an equation for remagnetization dynamics. This equation is treated here. Several processes occurring at different timescales play a role in the description
of these dynamics. We use a phenomenological equation for the dynamics of laser-induced remagnetization, derived by Józsa [25]. The different contributions in this equation will be discussed in the following sections.

**Sub-picosecond demagnetization**

According to Beaurepaire et al. [9] (section 1.3) a fast demagnetization will take place within about 100 fs after laser excitation. In the phenomenological description of Józsa [25] the sub-picosecond demagnetization is not taken into account. The starting point of the description is the point at which the fast demagnetization is just finished.

The demagnetization is followed by a remagnetization which takes a couple of picoseconds. This fast remagnetization is caused by electron-phonon interactions involving hot electrons relaxing due to energy dissipation to the lattice [9]. This process can be described by an exponential decay:

$$\Delta M(t) = A_1 e^{-t/t_{ep}},$$

(2.19)

where $A_1$ is the amplitude before remagnetization, $t_{ep}$ is the electron-phonon relaxation time (typically 400 fs) and $\Delta M(t)$ is the change of the magnetization.

**Heat flow to the substrate**

After the electron-phonon relaxation in the first picoseconds the thin metal film will transfer heat to the substrate. The ease at which this heat can flow determines the evolution of the temperature distribution and the remagnetization in time. Two cases can be distinguished [25]:

- Heat can flow easily into the substrate. This causes a dynamic temperature profile which quickly reaches equilibrium (typically 100 ps). The relaxation of the temperature in the magnetic layer - and therefore the remagnetization - will in this case be governed by an exponential behavior.

- Heat flow to the substrate is predominantly blocked by an insulating layer. In this case the insulating layer is slowly heated by diffusion. Since the heat slowly flows out of the sample a nearly static temperature profile is present in the sample. The temperature gradient in the insulator is Gaussian and the remagnetization in the magnetic layer is in this case governed by a $1/\sqrt{t}$ dependence.
The samples used in our experiments contain a thin insulating oxidized layer in between the substrate and the platinum. Therefore, the remagnetization is, in practice, expected to be described by a combination of both cases. Implementing both behaviors into the remagnetization formula 2.19 yields

$$\Delta M(t) = A_1 e^{t/t_{\text{ep}}} + \frac{A_2}{\sqrt{t + t_0}} + A_3 e^{t/t_b}, \quad (2.20)$$

in which $A_2$ and $A_3$ are again amplitudes and $t_0$ is the time it takes for a Gaussian temperature profile in the insulating oxide layer to set in. The timescale $t_b$ describing the heat flow to the substrate is much larger than the electron-phonon relaxation time $t_{\text{ep}}$.

**Damped precessional motion**

The final term is due to a damped precessional motion governed by the LLG equation 2.9 (see section 2.2.1). This term can be written as the product of an exponential decay (the damping) and an oscillation. Thus, the complete equation for the remagnetization in a magnetic layer after excitation by a laser pulse becomes $[25]$

$$\Delta M(t) = A_1 e^{t/t_{\text{ep}}} + \frac{A_2}{\sqrt{t + t_0}} + A_3 e^{t/t_b} + A_4 e^{t/t_d} \sin(\omega t + \phi_0), \quad (2.21)$$

with $t_d$ the precessional decay time (much larger than $t_{\text{ep}}$), $\omega$ the frequency of the precession and $\phi_0$ the initial phase. In the research presented in this thesis it is expected that precessions in both the in-plane and the out-of-plane layer will be measured. Therefore, equation 2.21 is modified by adding a second precession:

$$\Delta M(t) = A_1 e^{t/t_{\text{ep}}} + \frac{A_2}{\sqrt{t + t_0}} + A_3 e^{t/t_b} + A_4 e^{t/t_d,\text{oop}} \sin(\omega_{\text{oop}} t + \phi_{0,\text{oop}})$$

$$\quad + A_5 e^{t/t_d,\text{ip}} \sin(\omega_{\text{ip}} t + \phi_{0,\text{ip}}). \quad (2.22)$$

The subscripts 'oop' and 'ip' indicate the precessions of the magnetization of respectively the out-of-plane and the in-plane magnetized layer.

In our research, the component of the magnetization, perpendicular to the film plane is measured. We define this component as $M_z$. Figure 2.9 shows a typical remagnetization curve as described by equation 2.22 for a non-collinear magnetic bilayer, which is depicted at the top right. $M_z$ is plotted as a function of the delay time $t$. The fast demagnetization is just finished at $t = 0$. 
CHAPTER 2. THEORY

Figure 2.9: A typical laser-induced remagnetization curve for a non-collinear magnetic bilayer. The perpendicular component of the magnetization is plotted as a function of the delay time. This component is also measured in our experiments. The fast demagnetization is just finished at $t = 0$.

2.3 Mechanisms for inducing a spin current

In this thesis, we aim on the demonstration of spin-polarized currents, induced by a femtosecond laser pulse. The generation of such currents can be explained by two separate excitation mechanisms. First, we present a mechanism called superdiffusive transport, which is an excitation occurring at sub-picosecond timescales. Next, the spin-dependent Seebeck effect is treated.

2.3.1 Superdiffusive transport

Superdiffusive transport is a type of transport in which electrons do not travel collisionless (ballistic), but also do not collide very often (diffusive). We will now explain the generation of a superdiffusive current due to femtosecond laser excitation on both the majority and minority spins.

Let us take a look at the schematic density of states of a ferromagnet in Figure 2.10a. We see that the density of states around the Fermi energy
$E_F$ differs for majority (up) and minority (down) spins. Now, the energy of the laser pulse is primarily absorbed by the electron system. The electrons around the Fermi level are excited to energy levels of about 1 eV [26]. The system is now in a strong nonequilibrium state (Figure 2.10b). It can be seen in Figure 2.10c that there are a lot more empty states for the minority spins to relax to than for the majority spins. This is illustrated by the green and the red boxes in Figure 2.10c for available states around the Fermi level. Consequently, majority and minority spins have different lifetimes. Zhukov et al. [26] obtained lifetimes from 5 fs (3 eV) up to 40 fs (0.5 eV) for the majority spins in Ni. The minority spins were found to have a spin-independent lifetime of only 2 fs.

At first sight, the lifetime of the majority spins seems too small for the electron to travel over significant distances. Still, due to their extremely high velocities of about 1 nm/fs ($\sim c/300$), the excited spins can travel over distances of tens of nanometers before they scatter [14]. In this way, a superdiffusive spin-polarized current of majority spins will flow out of the magnetic material, which was already shown in Figure 1.6 in section 1.3. As this cur-
rent is spin-polarized, it is accompanied by a demagnetization of the magnetic layer. Battiato et al. [14] predict that the outflow of spins provides a considerable contribution to the laser-induced sub-picosecond demagnetization [9] (see subsection 2.2.4) and can even completely explain it. So concluding, superdiffusive transport could be a possible excitation mechanism of spin-polarized currents.

2.3.2 Spin-dependent Seebeck effect

Another possible mechanism for exciting a spin-polarized current is the spin-dependent Seebeck effect. This effect is the phenomenon that a spin current is induced due to diffusion when a temperature gradient is present in a ferromagnetic material [27]. Unlike superdiffusive transport, the Seebeck effect can exist for relatively long timescales, since the current keeps flowing as long as a temperature gradient is present.

The general Seebeck effect is the effect, that an electric voltage $V$ is induced in a conductor when it is placed within a temperature difference $\Delta T$. The Seebeck effect is governed by the simple equation [27]

$$V = S \Delta T,$$

in which $S$ is the Seebeck coefficient. If a magnetic material is placed in a temperature gradient a similar effect occurs. However, in this case, $S$ differs for majority and minority spins, since their lifetimes are different (see 2.3.1).
Therefore, the gradient-induced current will be spin-polarized. This is called
the spin-dependent Seebeck effect.

In this thesis, magnetic bilayers are excited by a laser pulse. A possible
temperature distribution is schematically drawn in Figure 2.11 for an arbi-
trary moment after laser excitation. The temperature gradients induce a
spin-polarized current, according to the spin-dependent Seebeck effect.

We have seen that spin-polarized currents can either be induced by su-
perdiffusive transport or the spin-dependent Seebeck effect. Superdiffusive
transport induces a spin current by the outflow of high-energetic majority
spins, excited by the laser pulse. The spin-dependent Seebeck effect induces
a spin current due to temperature gradients in the sample, induced by laser
heating.
CHAPTER 2. THEORY
Chapter 3

Experimental

In this chapter we will introduce the experimental setup with which the time-resolved magnetization dynamics are measured after pulsed laser excitation. We start by discussing the physics of the magneto-optical Kerr effect (MOKE), which is an optical effect that is utilized to measure the magnetization of thin films. Second, we explain the time-resolved MOKE setup.

3.1 Phenomenological description of MOKE

The magneto-optical Kerr effect (MOKE) is named after John Kerr, who predicted that the polarization of a linearly polarized light beam rotates and becomes elliptic when reflecting on a magnetic surface [28]. The ellipticity and the rotation angle of the reflected laser beam are a measure for the magnetization of the material. The optical response of a material to an electric field is described by the dielectric tensor \( \varepsilon \). For a non-magnetic, optically isotropic material the off-diagonal elements of this dielectric tensor are zero. For a magnetized material, however, the dielectric tensor becomes

\[
\varepsilon = \begin{pmatrix}
\varepsilon_{xx} & \varepsilon_{xy} & \varepsilon_{xz} \\
-\varepsilon_{xy} & \varepsilon_{xx} & \varepsilon_{yz} \\
-\varepsilon_{xz} & -\varepsilon_{yz} & \varepsilon_{xx}
\end{pmatrix}.
\]  

(3.1)

The off-diagonal elements are now dependent on the magnetization \( \vec{M} \) of the material. In particular, element \( \varepsilon_{ij} \) depends on the component of \( \vec{M} \) in the direction parallel to \( i \times j \). Let us, for simplicity, consider a sample with a magnetization pointing in the z-direction. Equation 3.1 can in this case be reduced to

\[
\varepsilon = \begin{pmatrix}
\varepsilon_{xx} & \varepsilon_{xy} & 0 \\
-\varepsilon_{xy} & \varepsilon_{xx} & 0 \\
0 & 0 & \varepsilon_{xx}
\end{pmatrix}.
\]  

(3.2)
It turns out that two of the eigenvectors of the dielectric tensor are $e_+ = 1/\sqrt{2}(1,i,0)^T$ and $e_- = 1/\sqrt{2}(1,-i,0)^T$, representing left- and right-handed circularly polarized light. The superscript $T$ denotes a transposed vector. The third eigenvector is $e_z = (0,0,1)^T$ which is just the unit vector in the propagation direction $z$. Diagonalizing the dielectric tensor above will result in a tensor which is mapped on these eigenvectors. The elements of this tensor are a measure for the magneto-optical response of light beams polarized in the direction of the corresponding eigenvector. Let us define the new basis by $P = (e_+^T, e_-^T, e_z^T)$. Diagonalization of equation 3.2 then yields

$$
\bar{\epsilon}_D = P^{-1}\bar{\epsilon}P = \begin{pmatrix}
\epsilon_{xx} + i\epsilon_{xy} & 0 & 0 \\
0 & \epsilon_{xx} - i\epsilon_{xy} & 0 \\
0 & 0 & \epsilon_{xx}
\end{pmatrix}.
$$

(3.3)

It can be seen that the elements - and hence the magneto-optical response - for left- and right-handed circularly polarized light differ: $\epsilon_\pm = \epsilon_{xx} \pm i\epsilon_{xy}$. This inequality is the mathematical origin of the Kerr effect. Namely, the incident linearly polarized light beam can also be seen as a superposition of two circularly polarized light beams: one left-handed, the other right-handed. Since the optical response of the material differs for both directions the reflected beam rotates and becomes elliptically polarized. This phenomenon is illustrated in Figure 3.1. Thus, linearly polarized light experiences a complex rotation $\psi = \psi_r + i\psi_e$ when reflecting on a magnetic surface, where $\psi_r$ and

![Figure 3.1: The MO Kerr effect. Linearly polarized light is a superposition of left- and right-handed circularly polarized light. The magneto-optical response is different for both polarizations. Therefore, linearly polarized light experiences a complex rotation upon reflection from a magnetic medium. Courtesy to Dalla Longa [29].](image)
\( \psi_e \) are respectively the induced magneto-optical rotation and ellipticity. The complex rotation is linked to the magnetization of the medium by [21]

\[ \psi = FM, \quad (3.4) \]

with \( F \) a generalized Fresnel coefficient.

Equation 3.4 can be generalized for time-resolved measurements of the magnetization in a perturbed system. Assuming relatively small changes in the Kerr rotation \( \psi \), equation 3.4 can be written as [30]

\[ \Delta \psi(t) = M_0 \Delta F(t) + F_0 \Delta M(t), \quad (3.5) \]

where \( M_0 \) and \( F_0 \) are unperturbed values and \( \Delta \) denotes perturbations. On the relevant timescales of this research - \( \sim 100 \) ps - the Fresnel coefficient \( F \) can assumed to be constant [30]. Therefore, \( \Delta \psi(t) \) and \( M(t) \) are proportional:

\[ \Delta \psi(t) \approx F_0 \Delta M(t). \quad (3.6) \]

This means that by measuring the Kerr signal as a function of time we can determine the temporal evolution of the magnetization. In the next section we will describe the setup used for such time-resolved Kerr measurements.

### 3.2 Time-Resolved Magneto-Optical Kerr Effect

In this thesis we are interested in the way the magnetization of a magnetic sample evolves in time after excitation by a laser pulse. Therefore, we use an all-optical pump-probe setup called Time-Resolved Magneto-Optical Kerr Effect setup or, in short, TR-MOKE setup. This setup exploits the Kerr effect to measure the change in magnetization, which we will discuss in detail below.

#### 3.2.1 Simple TR-MOKE setup

The simplest MOKE setup contains a laser, a beam splitter, two polarizers, a magnetic sample and a detector. We use a Tsunami (Spectra Physics) mode-locked Ti:Sapphire laser which generates pulses of approximately 70 fs at a frequency of 80 MHz. The output power is typically 650 mW, i.e. 8 nJ per pulse. An energy of 1 nJ is used for the excitation. The wavelength of the laser is approximately 790 nm.

The beam splitter divides the laser beam into a pump and a probe pulse [21]. A time-resolved signal is achieved by adding a delay-line to the path of
the probe pulse, thereby making the time interval between pump and probe excitation adjustable. This basic TR-MOKE setup is shown in Figure 3.2. The external magnetic field (depicted by the purple arrows) is optional. The different parts of the setup will now be discussed in more detail.

Figure 3.2: A basic pump-probe TR-MOKE setup. The desired signal is achieved by two polarizing the probe beam twice. Courtesy to Dalla Longa [29].

Separating the pump and probe pulse

The pump pulse (the thick line) excites the sample. Its intensity is about 95% of the intensity before the splitting. The weak probe pulse (5% intensity) hits the sample before or after the pump pulse, depending on the position of the delay-line. Both pulses should hit the sample at exactly the same spot for an accurate measurement. However, since only the probe pulse gives the magneto-optical response, the reflected pump pulse should not be detected. The separation of both pulses is achieved by dividing the objective lens into four quadrants. This is illustrated in Figure 3.3. Both pulses enter the objective lens at the top half and leave at the bottom half but their trajectories are mirrored. In this way the reflected pump pulse can be blocked while the probe pulse is send to the detector.
Figure 3.3: The separation of both pulses is achieved by dividing the objective lens into four quadrants.

Quantification of the MO response

The linear polarizers P and A are used for the measurement of the MO response of the probe pulse (see Figure 3.2). We make use of the Jones formalism [31] to quantify this response. Polarizer P linearly polarizes the probe pulse under an angle $\alpha_P$ relative to the vertical axis. The polarization afterwards can therefore be described by the normalized Jones vector $(\sin \alpha_P, \cos \alpha_P)^T$. Next, the probe pulse experiences a complex Kerr rotation upon reflection from the magnetic medium [28]. This rotation can be described by the Jones matrix [29]

$$S = r_s \begin{pmatrix} 1 & -\psi \\ \psi & r_p/r_s \end{pmatrix},$$

(3.7)

where $r_p$ and $r_s$ are complex reflection coefficients, independent of the magnetization. The vectors $\vec{s}$ and $\vec{p}$ are respectively the horizontal and vertical axes perpendicular to the propagation direction, see Figure 3.2. The magnetization-dependent part $\psi = \psi_r + i\psi_i$ is the complex Kerr rotation described in section 3.1. Thus, the resulting Jones vector after reflection is $S(\sin \alpha_P, \cos \alpha_P)^T$. Finally, the probe pulse propagates through the second polarizer, which has a polarization angle of $\alpha_A$ relative to the vertical axis. The intensity of the detected laser pulse is proportional to the squared inner
product between the polarization axis of P and the polarization direction of the incoming light and can thus be written as \[29\]

\[
I = R \left| (\sin \alpha_P, \cos \alpha_P) \mathcal{S} \left( \frac{\sin \alpha_A}{\cos \alpha_A} \right) \right|^2, \tag{3.8}
\]

with \( R = |r_s|^2 \) the reflectivity. When considering only the lowest orders of \( \alpha_A \) and \( \psi_r \), it can be derived that the change in the MO response due to the pump excitation is given by \[21\]

\[
\Delta I(t) = 2R_0 \alpha_A \Delta \psi_r(t) + \alpha_A^2 \Delta R(t), \tag{3.9}
\]

with \( R_0 \) the reflectivity before laser excitation and \( \Delta R(t) \) and \( \Delta \psi_r(t) \) the pump-induced changes at time \( t \) after excitation. Since the reflectivity \( R \) is not dependent on the magnetization of the sample, the reflectivity difference \( \Delta R(t) \) is an unwanted artifact that has to be kept as small as possible. Therefore, it can be seen from equation 3.9 that the polarization angle of the analyzer \( \alpha_A \) must be small.

### 3.2.2 Improved TR-MOKE setup

The sensitivity of the TR-MOKE setup can be improved by implementing a photo-elastic modulator (PEM), a chopper and two lock-ins. The resulting setup is shown in Figure 3.4. The functionality of these components will now be discussed. A photo-elastic modulator is a birefringent crystal that induces a phase difference between the left- and right-polarized components of the entering light. The PEM is put in the path of the probe pulse, behind the first polarizer P (see Figure 3.4). The polarization angle \( \alpha_P \) is set to 45°, yielding a Jones vector of \((1\sqrt{2}, 1\sqrt{2})^T\). The easy axis of the PEM is parallel to the s-axis, which is the horizontal axis perpendicular to the propagation (see Figure 3.4). In that case, the Jones matrix of the PEM is given by \[29\]

\[
M(t') = \begin{pmatrix} 1 & 0 \\ 0 & e^{A_0 \cos(\Omega t')} \end{pmatrix}, \tag{3.10}
\]

where \( A_0 \) is the amplitude of the retardation, \( \Omega \) is the operating frequency of 50 kHz and \( t' \) is a measure for the oscillations in the crystal referred to as the "internal time" by Dalla Longa \[29\]. The pump-induced change of the output signal \( \Delta I(t) \) can be divided into harmonics of the operating frequency of the PEM. For the specific orientation described above these harmonics can
be approximated by [21]

$$\frac{\Delta I_{1F}(t)}{I_{0f}} = 2J_1(A_0)\Delta \psi_s(t),$$  \hspace{1cm} (3.11a)$$
$$\frac{\Delta I_{2F}(t)}{I_{0f}} = 2J_2(A_0)\Delta \psi_r(t),$$  \hspace{1cm} (3.11b)$$

where $J_n(A_0)$ is the Bessel function of order $n$ at the retardation amplitude $A_0$ of the PEM. It follows from equations 3.11a and 3.11b that the signal of the first harmonic $F$ (50 kHz) is proportional to the Kerr ellipticity, while the second harmonic is proportional to the rotation.

The TR-MOKE can be further improved by adding a chopper in the path of the pump probe (see Figure 3.4). The chopper is a rotating disc that blocks the pump beam half of the time. We use a chopper which results in a 60 Hz modulation of the signal. The resulting signal reaching the detector is illustrated in 3.5. A double modulation is visible: a slow pump-induced change caused by the chopper and a fast oscillation due to the PEM. The measured signal is the pump-induced voltage difference indicated by $\delta V$. This signal is extracted by two lock-ins (L1 and L2 in Figure 3.4). The
detected signal is the input of L1, which is referenced with the PEM. Selecting the first or the second harmonic respectively yields the equilibrium Kerr ellipticity or rotation [21]. The output of L1 is then send to the input of L2, which is referenced with the chopper. According to equations 3.11a and 3.11b, this yields the time-resolved ellipticity or rotation (again dependent on the selected harmonic). Both variables are proportional to the time-resolved magnetization (equation 3.4). The sensitivity is now increased since all frequencies but the modulation frequency are filtered by the lock-ins.

Figure 3.5: Cartoon of the detected signal. The signal is modulated by both the chopper and the PEM. The difference indicated by $\delta V$ is the desired signal. Courtesy to Dalla Longa [29].
Chapter 4

Results

In this chapter, we demonstrate the presence of a laser-induced spin-transfer torque-induced precession of the magnetization in a non-collinear magnetic bilayer. First, we show that the magnetic layers of the bilayer are magnetostatically uncoupled and that the laser-induced precessions of the magnetization in both layers can be distinguished. Next, properties of these precessions, such as amplitude and phase, are analyzed with macro-spin models in order to find the excitation mechanisms that induced the precessions of the layers. It turns out that the in-plane magnetized layer is excited by a laser-induced spin-transfer torque, caused by a short spin-polarized current of majority spins, coming from the other magnetic layer. Then, the influence of the spacer layer on spin transport is investigated by substituting Cu with the strong spin scatterer Pt. Finally, possible microscopic origins of the measured spin currents are briefly discussed.

4.1 Static magnetic properties of the bilayer

Measurements of magnetization dynamics are performed on non-collinear magnetic bilayers in this research. We want to describe these dynamics for each magnetic layer separately. Therefore, it is important that these layers are magnetostatically uncoupled. The magnetostatic behavior of the non-collinear bilayer is investigated by measured hysteresis loops in this section.

In this research, measurements of magnetization dynamics are performed on Pt$_4$/Co$_{0.2}$/Ni$_{0.8}$/Co$_{0.2}$/X$_{0-20}$/Co$_3$/Pt$_1$ samples, in which Cu and Pt are used as the wedge-shaped spacer layer X. The thicknesses are given in nanometers. A cartoon of the sample is shown in Figure 4.1. The slope of the wedge is heavily exaggerated in the picture for illustrative purposes. The magnetizations of the layers in the sample are colored red and blue.
These colors will be associated with the bottom layer (red) and the top layer (blue) in the remainder of this thesis. The gray middle layer is a non-magnetic spacer layer. We define the out-of-plane axis as the z-axis. We are primarily sensitive to changes of the magnetization in this direction, which is a consequence of the used MOKE setup.

The Co/Ni multilayer is out-of-plane magnetized due to the interaction at the Co/Ni interface and interaction with the Pt layer underneath. The thick Co layer on top is magnetized in-plane. The top Pt layer prevents the sample from oxidation. The samples are deposited on a B-doped Si substrate which is a good heat conductor. The heat of the laser pulse can thus flow out of the sample easily. The samples are grown with DC magnetron sputtering.

The magnetic layers in our sample need to be magnetostatically uncoupled for a separate description of the magnetization dynamics to be valid. To investigate the role of uncoupling, static hysteresis loops are obtained. Before presenting the results of the measurements, we show what a typical hysteresis loop would look like if the layers are indeed uncoupled.

In section 2.1.3, we already showed typical hysteresis loops for in-plane and out-of-plane magnetized layers in the case of an applied field, perpendicular to the film plane (Figure 2.5). The bilayer in Figure 4.1 contains both of these layers. So, provided that the layers are uncoupled, a hysteresis measurement of the bilayer should yield a superposition of the hysteresis loops.
of the separate layers. A typical hysteresis loop is schematically depicted in Figure 4.2. The z-component of the total magnetization of the sample, $M_z$, is plotted as a function of the applied field $B$, which is applied in the z-direction. It can be seen that the in-plane layer is tilted slightly out-of-plane by the applied field. Also, the out-of-plane layer switches when the applied field equals the coercive field $B_c$ (see section 2.1.3).

![Figure 4.2: Schematic hysteresis loop of a non-collinear uncoupled bilayer. The orientation of the applied field is shown on the top left. The perpendicular component of the total magnetization is plotted as a function of the applied field. The colored arrows indicate the orientation of the magnetization in both layers. The black arrows indicate the sweep direction.](image)

We thus expect to measure the hysteresis loop of Figure 4.2 if there is no coupling. The experimental results will now be presented. Hysteresis loops are obtained for increasing spacer layer thicknesses. The results are shown in Figure 4.3. In Figure 4.3a, coercive fields, deduced from the hysteresis loops, are plotted versus the spacer layer thickness $d$. Both layers appear to be coupled until $d \approx 2$ nm. After that, the coercive field is fixed and the bottom layer seems to act as an isolated magnet. Figure 4.3b shows some examples of corresponding hysteresis loops. The Roman numbers correspond to the red data points in Figure 4.3a. The layers are clearly coupled without a spacer layer (0 nm), as both layers act as a single in-plane layer. Since there is no visible switch of the bottom layer in this case, a coercive field can not be defined. Hence, there is no corresponding data point for $d = 0$ nm in Figure 4.3a.

From a spacer layer thickness of 2 nm, we see the typical hysteresis as predicted in Figure 4.2: the top layer is gradually tilted out-of-plane by the
applied field while the bottom layer suddenly switches at the coercivity field (about 20 mT). This verifies that the role of coupling appears to be small for spacer layer thicknesses larger than 2 nm.

The time-resolved measurements were only performed at spacer layers thicknesses for which both layers seem to be uncoupled. Although a minor contribution of dipole or exchange coupling could still be present in our time-resolved measurements, we assume complete uncoupling from now on.

We have seen that switching of the bottom layer can be achieved by the application of an opposing, perpendicularly oriented magnetic field. On similar grounds, the in-plane magnetized top layer can be switched by applying a magnetic field which is directed parallel to the film plane. Consequently, the possible orientations of the magnetization of both layers - \( \rightarrow \uparrow, \rightarrow \downarrow, \leftarrow \uparrow \) and \( \leftarrow \downarrow \) - can be chosen in advance of every measurement.

![Figure 4.3](image)

**Figure 4.3:** a) Measured coercive fields for increasing spacer layer thickness \( d \). Uncoupling of the layers appears at \( d \approx 2 \) nm and seems to be maintained for \( d > 2 \) nm as the coercivity stays the same. b) Examples of obtained hysteresis loops at varying thicknesses. The Roman numbers correspond to the red data points in Figure a. For \( d = 0 \) nm, there is no coercivity and, therefore, no corresponding data point in Figure a. The loops indeed show uncoupling for \( d > 2.4 \) nm.
Concluding, it is shown in this section that the magnetic layers of our samples appear to be uncoupled if the spacer layer is thick enough. The role of dipole and exchange coupling is therefore assumed to be negligible in the forthcoming measurements. Furthermore, we have shown that the orientation of the individual magnetization of both layers - the configuration - can be set in advance of a time-resolved measurement by applying a magnetic field.

4.2 Acquisition of precessional properties

In this research, the mechanisms inducing a precession after pulsed laser excitation are deduced from the properties of these precessions, such as the frequency, amplitude and phase. The mechanisms that induced the precessions will eventually be revealed by comparing these precessional parameters to the macro-spin calculations described in section 2.2.2. It is explained in this section how these properties are obtained from TR-MOKE measurements.

![Graph showing precessional properties](image)

**Figure 4.4:** Top: typical measurement of the time-resolved Kerr rotation, performed on the sample depicted at the top right for the configuration $\alpha_0$. The obtained remagnetization curve is fitted to equation 2.22. Bottom: precessional parts of the time-resolved Kerr rotation for both layers, isolated from the fit of the top graph. The different precessional properties are indicated.
The (damped) precessions of the magnetization in each magnetic layer can be described by

\[ \Delta M_z(t) = A \sin(2\pi ft + \phi), \]  

(4.1)

in which \( A \), \( f \) and \( \phi \) are respectively the amplitude, frequency and phase of the precessions. We want to deduce these precessional properties from our TR-MOKE measurements. The procedure is presented next.

As an example, we show a typical TR-MOKE measurement, performed on the sample of Figure 4.1 for a 4 nm Cu spacer layer. A magnetic field of 70 mT is applied in-plane. The configuration of both layers is \( \uparrow \downarrow \). The measurement is shown in Figure 4.4.

The top graph shows the time-resolved Kerr rotation after pulsed laser excitation, which is proportional to the change of the z-component of the magnetization \( \Delta M_z(t) \), according to equation 3.6. The time \( t \) is the delay between the pulsed laser excitation and the arrival of the probe pulse, i.e. the sample is excited by the laser at \( t = 0 \). In section 2.2.4, the remagnetization after laser excitation is described by the phenomenological equation 2.22. The top graph of Figure 4.4 is fitted to this equation (the green line), yielding values for all 14 variables in equation 2.22 and, in particular, for the precessional properties of both layers. We developed a Graphical User Interface (GUI) in Matlab for fitting the measurements.

Now that the values of \( A \), \( f \) and \( \phi \) are known, the precessions of both layers can be isolated from the fit. The obtained precessions, governed by equation 4.1, are plotted in the bottom graph of Figure 4.6, where also the different properties are indicated.

To conclude, it is shown in this section that the precessional properties of both layers of the bilayer can be obtained by fitting a TR-MOKE measurement to the phenomenological equation 2.22. In the forthcoming sections, these properties are used to demonstrate the presence of a spin-transfer torque in the magnetic bilayer.

### 4.3 Distinguishing the magnetic layers

It is demonstrated in the previous section how the precessions of the magnetization in both layers can be isolated from a TR-MOKE measurement. For each layer, we want to separately compare the properties of these precessions to macro-spin models in order to find out which mechanisms excite the layers. However, from the data in Figure 4.4 we cannot assign the observed precessions to a specific layer. Consequently, before we can compare the precessional dynamics to macro-spin models, the precessions have to be separated, which is the main goal of the results presented in this section.
The top layer of our samples is magnetized in-plane and the bottom layer is magnetized out-of-plane (Figure 4.1). The precession of the magnetization, described by the LLG equation 2.9 will therefore be governed by different dynamics in each layer. In particular, the frequencies of precessions in both layers show a completely different magnetic field-dependence. It was shown in section 2.2.2 that, for small applied fields, the frequencies in the in-plane layer $f_{ip}$ and the out-of-plane layer $f_{oop}$ are respectively given by:

$$f_{ip} = \frac{\gamma}{2\pi} \sqrt{B_{app}(B_{app} + \mu_0 M_{sat})},$$

$$f_{oop} = \frac{\gamma}{2\pi} \sqrt{(B_{ani}^2 - B_{app}^2)},$$

with $B_{app}$ the in-plane applied field in T, $\gamma$ the gyromagnetic ratio, $\mu_0$ the vacuum permeability, $M_{sat}$ the saturation magnetization of the in-plane layer and $B_{ani}$ the anisotropy field in the out-of-plane layer in T.

This means that the layers can be distinguished by a measurement of the frequency as a function of the applied field strength. We performed such a measurement on the bilayer of Figure 4.1 for a Cu spacer layer thickness of 4 nm. A magnetic field is applied in-plane. Its strength is varied between -125 and 125 mT. The used configuration is thus $\parallel \downarrow$. For each applied field the frequencies of the laser-induced precessions of both layers have been obtained from the raw data, following the procedure described in section 4.2. For $B_{app} \approx 50$ mT both frequencies are almost equal and therefore hard to separate. For even smaller fields, a precession should only be measured in the out-of-plane layer ($f_{ip}$ goes to 0, see equation 4.2). However, as such a precession damps out within the first oscillation, measuring its frequency is difficult. Therefore, there are no data points for $B_{app} < 50$ mT.

The final results are depicted in Figure 4.5. It can be seen that both layers can indeed be easily distinguished. The data points of the top layer are fitted to equation 4.2, yielding a value of $1.1 \cdot 10^6$ A/m for $M_{sat}$ in the in-plane Co layer. The value is slightly smaller with respect to bulk Co. This might be caused by a contribution of an out-of-plane surface anisotropy due to the top Pt layer. The bottom layer is fitted to equation 4.3. We find a value of 250 mT for the perpendicular anisotropy field $B_{ani}$ in the out-of-plane Co/Ni layer.

So concluding, we have seen that the excitations of both magnetic layers can be distinguished by a field-dependent frequency measurement. In the following sections, the corresponding precessional phases and amplitudes are analyzed to find out if laser-induced spin-transfer torques caused the obtained precessions.
CHAPTER 4. RESULTS

4.4 Excitation mechanisms of laser-induced precessions

The goal of this research is to demonstrate that precessions of the magnetization of the magnetic layers of a non-collinear bilayer are caused by a laser-induced spin-transfer torque. However, apart from a spin-transfer torque, there are other excitation mechanisms that can induce a precession: an anisotropy pulse (a laser-induced change of the anisotropy), dipole coupling between the layers and exchange coupling between the layers. We assume that excitations, induced by coupling, are unimportant, since our measurements are only performed at spacer layers thicknesses, for which the layers are determined to be uncoupled (section 4.1). Therefore, we only need to consider the possibility of spin-transfer torque-induced and anisotropy pulse-induced excitations of the layers. In this section, it is shown that we can distinguish these two mechanisms by looking at the phases of the laser-induced precessions.

Figure 4.5: Measured frequencies of the precessions of the magnetization in the top layer (blue) and the bottom layer (red) as a function of the in-plane applied field. The top layer and the bottom layer are fitted by respectively equation 4.2 and 4.3.
4.4.1 Phase measurement

Again, we consider the in-plane field-dependent TR-MOKE measurements of the previous section. We focus on the obtained phases, corresponding to the frequencies of Figure 4.5. The phases are plotted in Figure 4.6 as a function of the applied field. The configuration of both layers is again \( \uparrow \downarrow \). It can be seen that the phases of both precessions, defined by equation 4.1, have a field-independent value of approximately \(-\pi/2\).

In the following, macro-spin models are analyzed to determine which phase an induced precession would have for all the possible excitation mechanisms. If this phase is not equal to \(-\pi/2\), the obtained precessions could not have been initiated by the considered mechanism. Within the two main mechanisms - spin-transfer torque and anisotropy pulse - further distinctions can be made. Spin-transfer torque can be induced by majority or minority spin currents and anisotropy pulse can cause an increase or a decrease of the anisotropy. Furthermore, all of the excitations can be either long or short with respect to the time of one evolution of a precession. All possible excitation mechanisms are listed in Figure 4.7.

![Figure 4.6: Measured phases of the precessions of the magnetization in the top layer (blue) and the bottom layer (red) as a function of the in-plane applied field. The purple line is \( \phi = -\pi/2 \).](image-url)
This makes a total of 8 possible excitation mechanisms for each layer, for which we have to compare the theoretical phase to the obtained value of $-\pi/2$. This is quite a laborious job. However, all 16 (two layers) analyses are performed in a similar way. One example of how the theoretical phases are obtained can be found in section 2.2.3. In the next section, the obtained phase is compared to theoretical phases for different excitation mechanisms.

4.4.2 Comparing phases: possible excitation mechanisms

In this section, we present the theoretical phases for precessions induced by all possible excitation mechanisms. These phases are then compared to the measured phase of $-\pi/2$, obtained in the previous section. In this way, we determine which mechanisms could have possibly excited the magnetic layers in our sample.

The example of 2.2.3 shows how the phases are obtained. The remaining 14 excitation mechanism from the list of section 4.4.1 are treated in a similar fashion. The obtained phases are depicted in Figure 4.7. Only mechanisms, for which the phase matches the measured value of $-\pi/2$ (section 4.4.1), could have excited the layers.

There are no anisotropy pulse-induced precessions possible in the in-plane layer. This is a consequence of applying the magnetic field parallel to the film plane. The magnetization and the anisotropy field initially points along the applied field. Therefore, a change in the anisotropy field will only change the size of the effective field, but not its direction and thus no precession can be induced. Consequently, the measured excitation of the top layer could have only been induced by a laser-induced spin-polarized current. Comparing the value of $-\pi/2$ (section 4.4.1) to the phases of Figure 4.7, it can be concluded that this current must be short and consisting of majority spins ($\uparrow$-spins for this configuration).

The out-of-plane bottom layer can, in principle, be excited by all 8 excitation mechanisms. Figure 4.7 shows that the phase measurement rules out 6 of these mechanisms: the bottom layer could have only been excited by a long-lasting decrease in perpendicular anisotropy or a short spin current of minority spins. Earlier experimental observations show that a minority spin current is unlikely [14, 26]. Therefore, we assume that an anisotropy pulse excited the bottom layer.
The measured precessions could have only been induced by mechanisms, for which the phase matches the measured value of $-\pi/2$. Thus, the top layer must have been excited by a short spin-transfer torque, caused by a laser-induced current of majority spins.

To conclude, we have demonstrated that the top layer is excited by a short majority spin-transfer torque upon pulsed laser excitation by analyzing the phases of the measured precessions. In the next section, the corresponding field-dependent amplitudes will be compared to macro-spin models to verify this conclusion.

### 4.5 Field-dependent amplitude

In this section, we focus on the measured precessional amplitudes as a function of applied field, and compare the results to macro-spin calculations. The amplitude of an induced precession depends differently on the in-plane applied field strength for every excitation mechanism. It is shown in the following that the measured field-dependent amplitudes are in agreement with a spin-transfer torque in the top layer and an anisotropy pulse in the bottom layer.
Figure 4.8: Measured amplitudes of the precessions of the magnetization in the top layer (blue) and the bottom layer (red) as a function of the in-plane applied field. The top layer is fitted by equation 4.4 for a spin-transfer torque-induced excitation and the bottom layer is fitted by equation 4.5 for an anisotropy pulse-induced excitation. Both fits are in good agreement with the data.

For both layers, macro-spin calculations are performed to obtain expressions for the precessional amplitude as a function of the applied field. An example of such a macro-spin model can be found in section 2.2.2. In these models, the top layer is excited by a short spin-transfer torque pulse and the bottom layer is excited by a long anisotropy pulse. A full derivation of the expressions can be found in Appendices B.2 and A.3. The field-dependent amplitudes in the top layer and the bottom layer are respectively given by

$$A_{z,\text{ip}}^{\text{STT}} \approx C_{\text{ip}} \left(1 - \frac{B_{\text{app}}}{\mu_0 M_{\text{sat}}} \right),$$

$$A_{z,\text{oop}}^{K} \approx C_{\text{oop}} B_{\text{app}}^2,$$

where $B_{\text{app}}$ is the applied field strength in T and $C_{\text{ip}}$ and $C_{\text{oop}}$ are constants to compensate for the arbitrary units of $A$. The superscripts $\text{STT}$ and $K$ indicate precessions induced by respectively a spin-transfer torque and an anisotropy pulse.

Figure 4.8 shows the amplitudes, associated to the frequencies and phases
The amplitudes of the in-plane layer are fitted to equation 4.4, in which we used the reported value of $1.1 \cdot 10^6$ A/m for $M_{\text{sat}}$ (section 4.3). The data shows a reasonable fit. The amplitudes of the out-of-plane layer are fitted to equation 4.5. The data shows the expected quadratic behavior.

It can be concluded that the measured field-dependent amplitudes are in agreement with a spin-transfer torque in the top layer and an anisotropy pulse in the bottom layer. It will be confirmed in the following section that a laser-induced spin-transfer torque excites the top layer by looking at the precessions as a function of the angle of the applied field.

### 4.6 Angle-dependent measurements

So far, our measurements are all performed with the external magnetic field applied along the film plane. In this section, precessions are measured for an applied field, which is tilted slightly out of the plane over varying angles. The angular dependencies of these precessions verify that the top layer is excited by a spin current.

We performed TR-MOKE measurements on the angular dependence of the amplitude and phase. In this case, a magnetic field is applied at an angle $\beta$ with respect to the film plane. As we are solely interested in precessions in the top layer, we do not want to induce any precessions in the bottom layer. This is achieved by applying a very small field of 25 mT, which will, according to the red curve in Figure 4.8, hardly excite the bottom layer.

The precessional properties are deduced from the raw measurement with the method described in section 4.2.

The amplitudes and phases are depicted in Figure 4.9 as a function of the angle $\beta$. The blue lines are guides to the eye. It turns out that the measured variables do not depend on $\beta$. Before we demonstrate that this is agreement with a laser-induced spin-transfer torque, we first show that the top layer is definitely not excited by an anisotropy pulse.

In our previous measurements, the applied field was directed along the film plane. The top layer could therefore only be excited by a spin current (see Figure 4.7). With the field now applied under an angle, the top layer can also be excited by an anisotropy pulse. The dashed red lines in Figure 4.9 schematically show what an angle-dependent measurement should look like if the top layer is induced by an anisotropy pulse. Figure 4.9a shows that the phase of the precessions should change sign when going from negative to positive angles. This is certainly not the case. Furthermore, the amplitude should depend heavily on $\beta$ and it must be zero if the field is applied in-plane,
CHAPTER 4. RESULTS

Figure 4.9: Measured frequencies, amplitudes and phases of the precessions of the magnetization in the top layer as a function of the angle $\beta$, at which the external field is applied. An angle of 0° denotes in-plane orientation. None of the measured variables depend on $\beta$. A macro-spin model (the horizontal lines) shows that these observations verify that a spin current is the mechanism that excites the top layer.

i.e. $\beta = 0^\circ$. This is in disagreement with the measured amplitudes. It can thus be concluded that the top layer is not excited by an anisotropy pulse.

We will now verify that the measured phases and amplitudes are in agreement with a spin-transfer torque-induced excitation in the top layer. In Appendix C the equilibrium angle $\alpha$ of the magnetization is calculated for a given applied field angle $\beta$ and an applied field of $B_{\text{app}} = 25$ mT. The angles $\alpha$ and $\beta$ are both defined with respect to the film plane. For the conditions of the measurement of Figure 4.9 an upper bound for the tilting angle of the magnetization is

$$\alpha < 0.2^\circ. \quad (4.6)$$

So, that means that, even for the largest value of $\beta$, the magnetization direction is not significantly influenced by the off-parallel applied field. This seems logical, since the anisotropy field that opposes the tilting of $\overline{M}$ is proportional to $M_{\text{sat}}$, which is over an order of magnitude larger than the applied field. The applied field is just not strong enough to compete with the demagnetization field, caused by the shape anisotropy in the magnetic
Besides the equilibrium direction of $\overrightarrow{M}$, the spin current is almost independent of $\beta$. Thus, we can expect identical precessions for all applied angles. Consequently, we can conclude that the frequency, amplitude and phase should not change as a function of $\beta$. This conclusion is confirmed by the results in Figure 4.9. Consequently, the statement that a spin current excites the top layer, is verified.

In conclusion, it is confirmed in this section that the top layer is excited by a laser-induced spin-transfer torque by studying the angular dependence of the precessional properties. In the following sections, the influence of the spacer layer material on the size of this spin-transfer torque is investigated.

### 4.7 Comparing precessions in copper and platinum spacer layers

We investigate the influence of the spacer layer material on the induced precessions in this section. When the spacer layer is a strong spin scatterer, no current-induced precession should be measured in the top layer. On the other hand, the bottom layer should still be excited by an anisotropy pulse. This hypothesis is verified in the following.

We performed TR-MOKE on the magnetic bilayer of Figure 4.1 for a 5 nm Pt spacer layer, instead of the Cu used before. Besides the use of a different spacer layer material, the sample is identical to the one used in the previous measurements. The spins of a majority spin current will scatter much more in the Pt than in the Cu spacer layer. Consequently, the amount of spin current reaching the opposing magnetic layer is expected to be much less in the Pt sample. Hence, the spin-transfer torque acting on the both layers should be smaller.

We proved that the in-plane layer of the Cu sample was excited by a spin-transfer torque. Hence, since the Pt spacer layer possibly decreases this torque significantly, we now expect only a minor precession of the top layer. For the bottom layer, the consequence of using Pt depends on the excitation mechanism. If the bottom layer was excited by a spin current in the Cu sample, we expect the precessions to be much smaller in this layer too. However, for an anisotropy pulse-induced precession, a possible spin current, coming from the top layer, does not play a role in the excitation and, consequently, nor does the spin conductance of the spacer layer. In this case, the precessions of the bottom layer should not depend on the spacer layer material.
We measured the phases and amplitudes of the precessions as a function of applied fields between 75 and 225 mT. An example of a TR-MOKE measurement is depicted in Figure 4.10 for an applied field of 200 mT. We see a clear difference in comparison to the identical measurement on the Cu spacer layer of Figure 4.4: only one precession can be detected in this case. We argued above that this is probably an excitation of the bottom layer. Again, we will verify this by looking at the precessional properties.

The obtained phases and amplitudes are shown in Figure 4.11. As expected, an excitation of the top layer is not detected. However, we did measure an excitation of the bottom layer. This must mean that the bottom layer is excited by an anisotropy pulse. The phases in Figure 4.11a are identical to the phases for a Cu spacer layer. The amplitudes of Figure 4.11b are fitted to 4.5. They show the quadratic field-dependence that we have seen earlier in Figure 4.8, which is typical for an anisotropy pulse.

So concluding, the fact that the excitations of the bottom layer are identical for a Cu and a Pt spacer layers confirms that the bottom layer is excited by a long anisotropy pulse. Furthermore, an excitation of the top
CHAPTER 4. RESULTS

Figure 4.11: Measured phases and amplitudes on the sample at the top left as a function of the applied field. A spin-transfer torque induced precession is not detected in the top layer. This can be expected, since the majority of a spin-current, coming from the bottom layer, will be scattered in the Pt before reaching the top layer. The excitation of the bottom layer must be induced by an anisotropy pulse. a) The phases are identical to the phase of the Cu sample. b) The amplitudes are fitted to 4.5 for an anisotropy pulse-induced precession. The data shows the expected quadratic behavior.

layer was not detected in the case of a Pt spacer layer, again demonstrating the laser-induced spin-transfer torque in the Cu sample. However, a minor current-induced excitation could still be present, though undetected due to the relatively large amplitudes for the anisotropy pulse mechanism. In the next section, the spin current exciting the top layer will be quantified as a function of spacer layer thickness for Cu and Pt.

4.8 Quantification of spin currents

In the previous section we concluded that the spin-transfer torque, acting on the in-plane layer must be small in the case of a Pt spacer layer. However, we did not detect a precession in the top layer at all. It is demonstrated in this section that a small precession is still induced in the top layer when using a Pt spacer layer, but that it is much smaller than the anisotropy pulse-
induced precessions in the bottom layer. Furthermore, the amplitudes of the precessions for Pt and Cu are compared in order to quantify the spin currents flowing through both materials as a function of spacer layer thickness.

In the previous section, a possible excitation of the top layer in the Pt sample could not be detected due to the major contribution of the excitation of the bottom layer to the Kerr rotation. So, in order to detect a precession of the top layer, the bottom layer should not be excited. This can be achieved by measuring without an applied field, according to equations 4.4 and 4.5. However, an unintentional remanent magnetic field of about 5 mT could still have been present during the measurements.

TR-MOKE measurements are performed on both the Cu and the Pt sample for varying spacer layer thicknesses $d$. Figure 4.12 shows an example at a thickness of 4 nm. Standard demagnetization curves, as described by Beaufrepaire et al. [9], are measured with an extra contribution of the spin current. These contributions are the deviations from the thermal background, illustrated by the red dashed lines. A clear contribution of a spin-transfer torque in the Cu sample is detected. More interesting is the fact that also a small contribution of spin transfer is visible in the Pt sample. Next, the torque-

![Figure 4.12: Measured Kerr rotations of a sample with a Cu spacer layer (left) and a Pt spacer layer (right). No external field is applied. The dashed red lines are guides to eye, indicating the Kerr rotations that would have been measured without the excitation of the top layers of both samples by a spin current. The contribution of a spin-transfer torque is larger for the Cu spacer layer, since Cu is the better spin conductor.](image)

induced contributions are isolated via a procedure described in Appendix D. Figure 4.13 shows the normalized Kerr rotation for the Cu and the Pt spacer layer, from which the exponential thermal parts are subtracted. It is clear that the amplitude of the induced precessions is significantly smaller for the Pt spacer layer. This was expected, as less spin current flows through the Pt in comparison to the Cu.

Finally, we will now quantify the spin current-induced amplitude of the magnetization. This is done by expressing the current-induced amplitude as
CHAPTER 4. RESULTS

Figure 4.13: Normalized contribution of spin-transfer torque to the total Kerr rotation. The data is obtained by subtracted the exponential, thermal part of the Kerr rotation from the raw data of Figure 4.12. A clear difference is visible between the Cu and the Pt spacer layer.

Figure 4.13: Normalized contribution of spin-transfer torque to the total Kerr rotation. The data is obtained by subtracting the exponential, thermal part of the Kerr rotation from the raw data of Figure 4.12. A clear difference is visible between the Cu and the Pt spacer layer.

a percentage of the thermal exponential part of the demagnetization curve. The procedure can be found in Appendix D. Percentages are shown for variable spacer layer thicknesses in Figure 4.14.

The data is fitted to the an exponential decay function $\Delta \theta_{\text{STT}} / \theta_{\text{Kerr,thermal}} = C \exp(-d/\lambda)$ with $C$ a constant, $d$ the spacer layer thickness and $\lambda$ the spin-transfer length. We obtain $\lambda_{\text{Pt}} = 1.6 \, \text{nm}$ and $\lambda_{\text{Cu}} = 18.0 \, \text{nm}$, which verifies that the spins indeed scatter more in the Pt layer.

We neglected the fact, that the sensitivity for the bottom layer decreases for increasing spacer layer thicknesses. The change of sensitivity could influence spin transfer length for both samples. Therefore, we expect that the inclusion of the optical sensitivity in the analysis will yield more accurate values for the spin-transfer lengths. However, the calculation of the optical sensitivity is beyond the scope of this thesis.

In this section it is concluded that replacing the Pt spacer layer by the Cu spacer layer significantly decreases the precessional amplitude of the top layer. This observation can be assigned to the smaller spin-transfer length in Pt, which is found to be ten times smaller than for Cu. Furthermore, it is
Figure 4.14: Quantification of the size of the spin current-induced excitations for the Cu (red) and the Pt (blue) spacer layer. The percentages are the torque-induced contributions to the Kerr rotation, normalized to the subpicosecond demagnetization. More information on the method of quantification can be found in Appendix D. The percentages are fitted to an exponential decay, yielding spin-transfer lengths of $\lambda_{\text{Pt}} = 1.6 \, \text{nm}$ for Pt and $\lambda_{\text{Pt}} = 18.0 \, \text{nm}$ for Cu. The large difference between both lengths shows that spin currents can indeed cover significantly longer distances in Cu than in Pt.

shown that the precessions of the bottom layer are not significantly influenced by the Pt spacer layer, demonstrating that these precessions are caused by an anisotropy pulse rather than spin-transfer torque.

### 4.9 Microscopic origins of the spin currents

We demonstrated the presence of spin-transfer torque-induced precessions in the in-plane magnetized top layer of non-collinear magnetic bilayers. These precessions are caused by majority spin currents with short lifetimes when compared to the typical time of one revolution of the magnetization. We will comment on the possible microscopic origins of these spin currents in the following.

In section 2.3 two excitation mechanisms were suggested: the spin-dependent Seebeck effect and superdiffusive transport. Spin-polarized currents with
short lifetimes can be generated by both excitation mechanism. This is obvi-
ous for superdiffusive transport, as this is a phenomenon that takes place at
sub-picosecond timescales (see section 2.3.1). The spin-dependent Seebeck
effect can also generate a short spin current, provided that the temperature
gradients in the sample vanish rapidly (see section 2.3.2). Unfortunately,
however, from our experimental data the two possible contributions cannot
be distinguished.

Confirmation on the microscopic origin of the spin currents can be achieved
by excluding one of both mechanism. For example, by calculating the tem-
perature profile in the sample just after laser excitation and then estimating
the theoretical size of spin currents, induced by the spin-dependent Seebeck
effect. This is, however, beyond the scope of this thesis.
Chapter 5

Conclusion and outlook

5.1 Conclusions

In this thesis we studied the precessional dynamics of a non-collinear magnetic bilayer after pulsed laser excitation. This bilayer consisted of an out-of-plane magnetized Co/Ni multilayer at the bottom, an in-plane magnetized Co top layer and a wedged non-magnetic Cu spacer layer. Analysis of the obtained hysteresis loops showed that the magnetic layers of this bilayer are magneto-statically uncoupled from a spacer layer thickness of 2 nm. A TR-MOKE setup is used to measure time-resolved remagnetization curves after exciting the sample with a femtosecond laser pulse. Properties of the precessions of the magnetization in both magnetic layers were then obtained from these curves. The precessions of both layers were distinguished by comparison of the obtained frequencies to developed macro-spin models.

By studying the phases of the precessions and comparing the results to a macro-spin model the origin of the induced precessions could be determined. The top layer is excited by a short spin current of majority spins and the bottom layer by a long-lasting decrease in perpendicular anisotropy. A macro-spin analysis of the obtained field-dependent amplitudes confirmed these conclusions. Further verification was obtained by measuring the precessional properties as a function of the applied field angle.

We substituted the Cu spacer layer for a Pt spacer layer - a strong spin scatterer - for our final measurements. Again, TR-MOKE measurements were performed, which showed that the Spin-transfer torque-induced precessions in the in-plane were heavily decreased in amplitude by the Pt spacer layer, as expected. Furthermore, the precessions of the out-of-plane layer remained virtually the same, as expected from an anisotropy pulse. Finally, the amplitudes of spin-transfer torque-induced precessions were compared
and quantified for both the Cu and the Pt spacer layer. We obtained a spin transfer length in Cu approximately ten times as large as the spin transfer length in Pt.

5.2 Outlook

The measured femtosecond laser-induced spin-transfer torque observed in the magnetic top layer seems promising for future spintronics devices relying on all-optical switching of magnetic bits. That all-optical switching is possible can be seen in experiments using higher laser fluences where substantial spin currents are created after pulsed laser excitation. By tuning the thickness of the layer that generates the spin current with respect to the switchable layer, all-optical switching should become possible. Next to the prospect of all-optical switching, the here presented experiments could be used as a tool to determine the amount of spin transport in ultrafast magnetization dynamics. By varying the magnetic materials and spacer layers the physics governing femtosecond transport can potentially be unraveled.
Appendix A

The out-of-plane layer

In the following section the properties of a precession in the out-of-plane magnetized Co/Ni bottom layer are analyzed. Expressions for the precessional frequency $\omega$ and the amplitude $A_z$ are derived in terms of the applied field strength. However, to this end, we have to define the initial conditions of the system before laser excitation. Therefore, we start by looking for the equilibrium states of the magnetization.

A.1 Equilibrium state before laser excitation

The equilibrated system is shown in Figure A.1. The applied field $H_{\text{app}}$ is lying in-plane. The anisotropy field, however, is pointing in the perpendicular direction, which is the easy axis of the magnetic layer. The anisotropy is given by

$$H_{\text{oop}} = \left( \frac{H_K}{M_{\text{sat}}} - 1 \right) M_z \hat{e}_z. \quad (A.1)$$

In equilibrium, the magnetization $\vec{M}$ points under an angle $\theta$ with respect to the vertical axis. The value of $\theta$ is not trivial, since $H_{\text{oop}}$ and $H_{\text{app}}$ are not parallel and the fact that $H_{\text{oop}}$ depends on $M_z$. An expression for this equilibrium angle is derived in the following. We make use of the LLG equation 2.9. The magnetization is given by

$$\vec{M}_{\text{eq}} = \begin{pmatrix} M_x \\ M_y \\ M_z \end{pmatrix} = \begin{pmatrix} 0 \\ M_{\text{sat}} \sin \theta \\ M_{\text{sat}} \cos \theta \end{pmatrix}. \quad (A.2)$$
Together with A.1 the cross product in 2.9 reads

\[ \vec{M}_{eq} \times \vec{H}_{eff} = \det \begin{pmatrix} \hat{e}_x & \hat{e}_y & \hat{e}_z \\ 0 & M_y & M_z \\ 0 & H_{app} & \left( \frac{H_K}{M_{sat}} - 1 \right) M_z \end{pmatrix}. \]  

(A.3)

As we are analyzing an equilibrated system, the left-hand side of 2.9 should be zero, i.e. \( \frac{d\vec{M}}{dt} = 0 \). A non-trivial equation is then only found in the x-direction:

\[ M_y \left( \frac{H_K}{M_{sat}} - 1 \right) M_z - M_z H_{app} = 0. \]  

(A.4)

Substitution of A.2 into A.4 yields

\[ M_{sat} \sin \theta \left( \frac{H_K}{M_{sat}} - 1 \right) M_{sat} \cos \theta - M_{sat} \cos \theta H_{app} = 0, \]  

(A.5)

\[ \Rightarrow \left[ (H_K - M_{sat}) \sin \theta - H_{app} \right] \cos \theta = 0. \]  

(A.6)

Equation A.6 has two solutions:

\[ \sin \theta_1 = \frac{H_{app}}{H_{ani}}, \]  

(A.7a)

\[ \cos \theta_2 = 0, \]  

(A.7b)

where \( H_{ani} = H_K - M_{sat} \). We will use this abbreviation in the remainder of this thesis. It will turn out later that condition A.7a is valid for \( H_{app} < H_{ani} \), whereas A.7b applies to larger fields.
A.2 Field-dependent frequency

In this section, the frequency of a precession in an out-of-plane magnetized layer is derived as a function of an in-plane applied field. The model is shown in Figure A.2.

The magnetization and the effective field are given by

\[
\mathbf{M} = \begin{pmatrix} M_x \\ M_y \\ M_z \end{pmatrix} = \begin{pmatrix} A_x \sin(\omega t) \\ -A_y \cos(\omega t) + M_{sat} \cos \theta \\ A_z \cos(\omega t) + M_{sat} \sin \theta \end{pmatrix}, \tag{A.8}
\]

\[
\mathbf{H}_{\text{eff}} = \begin{pmatrix} 0 \\ H_{\text{app}} \\ \frac{H_{\text{ani}}}{M_{\text{sat}}} M_z \end{pmatrix}, \tag{A.9}
\]

which leads to

\[
\mathbf{M} \times \mathbf{H}_{\text{eff}} = \det \begin{pmatrix} \mathbf{e}_x & \mathbf{e}_y & \mathbf{e}_z \\ M_x & M_y & M_z \\ 0 & H_{\text{app}} & \frac{H_{\text{ani}}}{M_{\text{sat}}} M_z \end{pmatrix} = \begin{pmatrix} M_y \frac{H_{\text{ani}}}{M_{\text{sat}}} M_z - M_z H_{\text{app}} \\ -M_x \frac{H_{\text{ani}}}{M_{\text{sat}}} M_y \\ M_x H_{\text{app}} \end{pmatrix}, \tag{A.10}
\]

and

\[
\frac{d \mathbf{M}}{dt} = \begin{pmatrix} A_x \omega \cos(\omega t) \\ A_y \omega \sin(\omega t) \\ -A_z \omega \sin(\omega t) \end{pmatrix}. \tag{A.11}
\]
By substituting A.8, A.10 and A.11 into the LLG equation 2.9 we obtain

\[ A_x \omega \cos(\omega t) = -\gamma \mu_0 \left( (-A_y \cos(\omega t) + M_{\text{sat}} \cos \theta) \frac{H_{\text{ani}}}{M_{\text{sat}}} - H_{\text{app}} \right) \times (A_x \cos(\omega t) + M_{\text{sat}} \cos \theta), \]  
\[ (A.12a) \]

\[ A_y \omega \sin(\omega t) = \gamma \mu_0 A_x \sin(\omega t) \frac{H_{\text{ani}}}{M_{\text{sat}}} (A_x \cos(\omega t) + M_{\text{sat}} \sin \theta), \]  
\[ (A.12b) \]

\[ A_z \omega \sin(\omega t) = \gamma \mu_0 A_x \sin(\omega t) H_{\text{app}}. \]  
\[ (A.12c) \]

Since every equation in this set has to hold for every arbitrary time \( t \), we are free to pick a time for which \( \sin(\omega t) \) or \( \cos(\omega t) \) equals 1. The set of equations can then be simplified to

\[ x \omega = \gamma \mu_0 [Ayz + By - Cz - D], \]  
\[ (A.13a) \]

\[ y \omega = \gamma \mu_0 Bx, \]  
\[ (A.13b) \]

\[ z \omega = \gamma \mu_0 Ex, \]  
\[ (A.13c) \]

in which \( x, y \) and \( z \) are abbreviations of the amplitudes \( A_x, A_y \) and \( A_z \). The capital letters are used to maintain a good overview on the calculation. The are given by

\[ A = \frac{H_{\text{ani}}}{M_{\text{sat}}}, \]  
\[ (A.14a) \]

\[ B = H_{\text{ani}} \cos \theta, \]  
\[ (A.14b) \]

\[ C = H_{\text{ani}} \sin \theta - H_{\text{app}}, \]  
\[ (A.14c) \]

\[ D = (H_{\text{ani}} \sin \theta - H_{\text{app}}) M_{\text{sat}} \cos \theta, \]  
\[ (A.14d) \]

\[ E = H_{\text{app}}. \]  
\[ (A.14e) \]

An expression for \( \omega \) in terms of the variables \( A \) to \( E \) can be found by solving A.13. We start with substituting A.13c into A.13a:

\[ x \omega = \gamma \mu_0 \left[ A \gamma \mu_0 E \frac{\gamma \mu_0 E x}{\omega} + By - C \frac{\gamma \mu_0 E x}{\omega} - D \right], \]  
\[ (A.15) \]

\[ x \omega = \gamma \mu_0 \left[ \gamma \mu_0 A E x y \frac{y \omega}{\omega} + By - \frac{\gamma \mu_0 C E x}{\omega} \frac{y \omega}{\gamma \mu_0 B} - D \right]. \]  
\[ (A.16) \]

Now, \( x \) is eliminated by implementing A.13b:

\[ \frac{y \omega}{\gamma \mu_0 B} \omega = \gamma \mu_0 \left[ \frac{\gamma \mu_0 A E y}{\omega} \frac{y \omega}{\gamma \mu_0 B} + By - \frac{\gamma \mu_0 C E}{\omega} \frac{y \omega}{\gamma \mu_0 B} - D \right], \]  
\[ (A.17) \]

\[ y \omega^2 = \gamma^2 \mu_0^2 \left[ A E y^2 + B^2 y - C E y - B D \right], \]  
\[ (A.18) \]

\[ \omega^2 = \gamma^2 \mu_0^2 \left[ A E y + B^2 - C E - \frac{B D}{y} \right]. \]  
\[ (A.19) \]
Next, the variables $A$ to $E$ are replaced again for the expressions in A.14 and the abbreviations $x$, $y$ and $z$ are discarded:

$$\frac{\omega^2}{\gamma^2 \mu_0^2} = \frac{H_{\text{ani}}}{M_{\text{sat}}} H_{\text{app}} A_y + H_{\text{ani}}^2 \cos \theta^2 - \left( H_{\text{ani}} \sin \theta - H_{\text{app}} \right)$$

$$\times \left[ H_{\text{app}} + \frac{H_{\text{ani}} M_{\text{sat}} \cos^2 \theta}{A_y} \right]. \quad (A.20)$$

In section A.1, two equilibrium states were derived from the equilibrium condition A.4: $\sin \theta_1 = \frac{H_{\text{app}}}{H_{\text{ani}}}$ and $\cos \theta_2 = 0$. First, we derive an expression for $\theta = \theta_1$:

$$\omega_1^2 = \gamma^2 \mu_0^2 \left[ \frac{H_{\text{ani}}}{M_{\text{sat}}} H_{\text{app}} A_y + H_{\text{ani}}^2 \cos \theta_1^2 \right]. \quad (A.21)$$

The $y$-component of the amplitude of the precession $A_y$ is assumed to be small. Therefore, the latter equation can be approximated by

$$\omega_1^2 = \gamma^2 \mu_0^2 H_{\text{ani}}^2 \cos \theta_1^2, \quad (A.22)$$

$$\omega_1 = \gamma \mu_0 H_{\text{ani}} \cos \theta_1. \quad (A.23)$$

Again implementing the equilibrium condition $\sin \theta_1 = \frac{H_{\text{app}}}{H_{\text{ani}}}$ yields

$$\omega_1 = \gamma \mu_0 H_{\text{ani}} \cos \left[ \text{arcsin} \left( \frac{H_{\text{app}}}{H_{\text{ani}}} \right) \right], \quad (A.24)$$

$$\omega_1 = \gamma \mu_0 H_{\text{ani}} \sqrt{1 - \left( \frac{H_{\text{app}}}{H_{\text{ani}}} \right)^2}, \quad (A.25)$$

$$\omega_1 = \gamma \mu_0 \sqrt{H_{\text{ani}}^2 - H_{\text{app}}^2}, \quad (A.26)$$

$$\omega_1 = \gamma \sqrt{B_{\text{ani}}^2 - B_{\text{app}}^2}, \quad (A.27)$$

with $B = \mu_0 H$. It can be seen that A.27 only holds for $B_{\text{app}} < B_{\text{ani}}$. Apparently, the solution for higher fields will be found with the second equilibrium angle, i.e. $\cos \theta_2 = 0$. Substitution into A.20 yields

$$\omega_2^2 = \gamma^2 \mu_0^2 \left[ \frac{H_{\text{ani}}}{M_{\text{sat}}} H_{\text{app}} A_y - \left( H_{\text{ani}} \sin \theta - H_{\text{app}} \right) H_{\text{app}} \right]. \quad (A.28)$$

The condition $\cos \theta_2 = 0$ demands that $\sin \theta_2 = 1$. For small $A_y$ we can then write

$$\omega_2^2 = \gamma^2 \mu_0^2 \left[ -(H_{\text{ani}} - H_{\text{app}}) H_{\text{app}} \right], \quad (A.29)$$

$$\omega_2 = \gamma \mu_0 \sqrt{(H_{\text{app}} - H_{\text{ani}}) H_{\text{app}}}, \quad (A.30)$$

$$\omega_2 = \gamma \sqrt{(B_{\text{app}} - B_{\text{ani}}) B_{\text{app}}}. \quad (A.31)$$
As expected, this solution is valid for $B_{\text{app}} > B_{\text{ani}}$. To summarize: the frequency of the precession in terms of the applied field $B_{\text{app}}$ is given by

$$\omega = \begin{cases} 
\gamma \sqrt{(B_{\text{ani}}^2 - B_{\text{app}}^2)} , & B_{\text{app}} < B_{\text{ani}} \\
\gamma \sqrt{(B_{\text{app}} - B_{\text{ani}})B_{\text{app}}} , & B_{\text{app}} > B_{\text{ani}} 
\end{cases} \quad (A.32)$$

This result is a special case of a more general expression for an out-of-plane film with an applied field in an arbitrary direction, derived by Layadi [32].

### A.3 Field-dependent amplitude of an anisotropy pulse-induced precession

The equilibrium conditions derived in the previous section are now used to derive an expression for the $z$-component of the amplitude $A_z$ of the precession induced by an anisotropy pulse. We use the model as illustrated in Figure A.3. In order to ensure that the used indexes are consistent with the remainder, we define the plane of movement as the $(y,z)$-plane. We start by choosing an initial anisotropy $H_{\text{ani}}$ and magnetization $\vec{M}_0$, of which the direction of the latter is given by $\theta_0$. This angle is given by A.7a (small applied fields are used). Next, we consider a small anisotropy pulse, yielding a slightly smaller anisotropy field $H_{\text{ani},1} = CH_{\text{ani}}$, where $C$ is a dimensionless number between 0 and 1. An effective field now arises, around which the magnetization $\vec{M}_1$ precesses until it points parallel to this field. $A_z$ is acquired by substracting the $z$-components of $\vec{M}_0$ and $\vec{M}_1$. The model is shown in Figure A.3. The initial and final magnetizations read

$$\vec{M}_0 = M_{\text{sat}}(\sin \theta_0, \cos \theta_0), \quad (A.33)$$
$$\vec{M}_1 = M_{\text{sat}}(\sin \theta_1, \cos \theta_1). \quad (A.34)$$

Therefore, $A_z$ is given by

$$A_z = M_{0,z} - M_{1,z} = M_{\text{sat}}(\cos \theta_0 - \cos \theta_1). \quad (A.35)$$
Using the equilibrium condition A.7a leads to

\[ A_z = M_{\text{sat}} \left[ \cos \left( \arcsin \left( \frac{H_{\text{app}}}{H_{\text{ani}}} \right) \right) - \cos \left( \arcsin \left( \frac{H_{\text{app}}}{H_{\text{ani},1}} \right) \right) \right], \quad (A.36) \]

\[ A_z = M_{\text{sat}} \left[ \sqrt{1 - \left( \frac{H_{\text{app}}}{H_{\text{ani}}} \right)^2} - \sqrt{1 - \left( \frac{H_{\text{app}}}{H_{\text{ani},1}} \right)^2} \right], \quad (A.37) \]

\[ A_z = M_{\text{sat}} \left[ \sqrt{1 - \left( \frac{H_{\text{app}}}{H_{\text{ani}}} \right)^2} - \sqrt{1 - \left( \frac{H_{\text{app}}}{CH_{\text{ani}}} \right)^2} \right]. \quad (A.38) \]

Taylor expanding the latter equation around \( H_{\text{app}} = 0 \) yields

\[ A_{z,\text{oop}} \approx \frac{(1 - C^2)M_{\text{sat}}}{2C^2B_{\text{ani}}^2} B_{\text{app}}^2 + O(B_{\text{app}}^4), \quad (A.39) \]

where \( B_{\text{app}} = \mu_0 H_{\text{app}} \) and \( B_{\text{ani}} = \mu_0 H_{\text{ani}} \) are both given in units of Tesla.

Equation A.39 shows that the amplitude depends heavily on the applied field. This can be explained by looking at the precession’s plane of movement. At first sight, A.39 implies that no precession is induced for small fields. This is not the case. Since the measurements are performed in the polar MOKE configuration, we are only sensitive to changes in the z-direction. So for small fields, the induced precessions are not detected, as they evolve in the (x,y)-plane. For increasing \( H_{\text{app}} \), the z-component of the precession becomes bigger and bigger as the magnetization is tilted gradually towards the in-plane hard axis.
Appendix B

The in-plane layer

In the following section the properties of a spin-transfer torque-induced precession in the in-plane magnetized Co top layer are analyzed. Expressions for the precession’s frequency $\omega$ and the amplitude $|M_z|$ are derived in terms of the applied field strength. We start with the frequency expression.

B.1 Field-dependent frequency

In this section, the frequency of a precession in an in-plane magnetized layer is derived as a function of an in-plane applied field. The model is shown in Figure B.1.

The magnetization and the effective field are given by

\[
\begin{align*}
\vec{M} &= \begin{pmatrix} M_x \\ M_y \\ M_z \end{pmatrix} = \begin{pmatrix} A_x \sin(\omega t) \\ M_{\text{sat}} \\ A_x \cos(\omega t) \end{pmatrix}, \\
\vec{H}_{\text{eff}} &= \begin{pmatrix} 0 \\ H_{\text{app}} \\ -M_z \end{pmatrix},
\end{align*}
\]

which leads to

\[
\begin{align*}
\vec{M} \times \vec{H}_{\text{eff}} &= \det \begin{pmatrix} \vec{e}_x & \vec{e}_y & \vec{e}_z \\ A_x \sin(\omega t) & M_{\text{sat}} & A_x \cos(\omega t) \\ 0 & H_{\text{app}} & -M_z \end{pmatrix} = \begin{pmatrix} -A_x \cos(\omega t)[M_{\text{sat}} + H_{\text{app}}] \\ A_x \sin(\omega t) A_x \cos(\omega t) \\ H_{\text{app}} A_x \sin(\omega t) \end{pmatrix},
\end{align*}
\]

and

\[
\frac{d\vec{M}}{dt} = \begin{pmatrix} A_x \omega \cos(\omega t) \\ 0 \\ -A_x \omega \sin(\omega t) \end{pmatrix}.
\]
Inserting B.3 and B.4 into the LLG equation 2.9 yields \( \alpha = 0 \)

\[
A_x \omega \cos(\omega t) = \gamma \mu_0 A_z \cos(\omega t) [M_{sat} + H_{app}], \tag{B.5a}
\]

\[
0 = -\gamma \mu_0 A_x \sin(\omega t) A_z \cos(\omega t), \tag{B.5b}
\]

\[-A_z \omega \sin(\omega t) = -\gamma \mu_0 H_{app} A_x \sin(\omega t). \tag{B.5c}
\]

These three equations should hold for every arbitrary time \( t \). Therefore, B.5a and B.5c can be written as

\[
A_x \omega = \gamma \mu_0 A_z [M_{sat} + H_{app}], \tag{B.6a}
\]

\[
A_z \omega = \gamma \mu_0 H_{app} A_x. \tag{B.6b}
\]

The desired expression is found by rewriting B.6b and substituting it into B.6a:

\[
A_x = \frac{A_x \omega}{\gamma \mu_0 H_{app}}, \tag{B.7}
\]

\[
\Rightarrow \frac{A_x \omega}{\gamma \mu_0 H_{app}} \omega = \gamma \mu_0 A_z [M_{sat} + H_{app}], \tag{B.8}
\]

\[
\omega = \gamma \mu_0 \sqrt{H_{app}(H_{app} + M_{sat})}. \tag{B.9}
\]

Finally, B.9 can be written as a function of the more convenient variable
\( B_{\text{app}} = \mu_0 H_{\text{app}} \), which is the magnetic field strength in Teslas:

\[
\omega = \gamma \mu_0 \sqrt{\frac{B_{\text{app}}}{\mu_0} \left( \frac{B_{\text{app}}}{\mu_0} + M_{\text{sat}} \right)}, \tag{B.10}
\]

\[
\omega = \gamma \sqrt{\mu_0 B_{\text{app}} \left( \frac{B_{\text{app}}}{\mu_0} + M_{\text{sat}} \right)}, \tag{B.11}
\]

\[
\omega = \gamma \sqrt{B_{\text{app}} \left( B_{\text{app}} + \mu_0 M_{\text{sat}} \right)}. \tag{B.12}
\]

## B.2 Field-dependent amplitude of a spin current-induced precession

In this section, the amplitude of a spin-transfer torque-induced precession in an in-plane magnetized layer is derived as a function of an in-plane applied field. The model is shown in Figure B.2. We first show the derivation and then an small-field approximation is presented.

### B.2.1 Derivation

The amplitude dependency in the in-plane layer is investigated by the model of Figure B.2. We consider the first picosecond after laser excitation: a spin-polarized current coming from the out-of-plane bottom layer causes a laser-induced spin-transfer torque in the in-plane layer. This torque is represented by a temporal field \( \vec{H}_T = H_T(\sin \alpha, \cos \alpha) \), of which the size is independent of \( H_{\text{app}} \). The direction of \( H_T \), however, does depend on \( H_{\text{app}} \): it is determined by the polarization of the incoming spins. We already proofed in section A.1 that the angle \( \alpha \) is given by the equilibrium condition A.7a (small applied fields are used). When the laser hits the sample the magnetization is almost instantly pulled parallel to the exchange field \( \vec{H}_{\text{ex}} \). The vertical component of the new magnetization \( \vec{M} \) determines the amplitude \( A_z \). The exchange field is given by

\[
\vec{H}_{\text{ex}} = \begin{pmatrix} 0 \\ H_T \sin \alpha + H_{\text{app}} \\ H_T \cos \alpha - M_z \end{pmatrix}. \tag{B.13}
\]

The direction of the equilibrium state of \( \vec{M} \) is governed by the equilibrated, undamped LLG equation 2.9:

\[-\gamma \mu_0 \vec{M} \times \vec{H}_{\text{ex}} = \frac{d\vec{M}}{dt} = 0. \tag{B.14}\]
This leads to

\[
\begin{pmatrix}
\mathbf{e}_x & \mathbf{e}_y & \mathbf{e}_z \\
0 & M_y & M_z \\
0 & H_T \sin \alpha + H_{\text{app}} & H_T \cos \alpha - M_z
\end{pmatrix}
= 0, \quad (B.15)
\]

\[
\Rightarrow M_y(H_T \cos \alpha - M_z) - M_z(H_T \sin \alpha + H_{\text{app}}) = 0, \quad (B.16)
\]

\[
\Rightarrow M_z = \frac{M_y H_T \cos \alpha}{M_y + H_{\text{app}} + H_T \sin \alpha}. \quad (B.17)
\]

In this case the z-component of the amplitude of the precession is equal to the z-component of \( \vec{M} \), i.e. \( A_z = M_z \). Furthermore, we expect that the magnetization is only slightly tilted out-of-plane. Therefore, we assume that \( M_y \approx M_{\text{sat}} \). We implement these assumptions and use the equilibrium condition A.7a once more:

\[
A_z = \frac{M_{\text{sat}} H_T \cos(\arcsin(H_{\text{app}} / H_{\text{ani,oop}}))}{M_{\text{sat}} + H_{\text{app}} + H_T H_{\text{ani,oop}}}, \quad (B.18)
\]

\[
A_z = \frac{M_{\text{sat}} H_T \sqrt{1 - (H_{\text{app}} / H_{\text{ani,oop}})^2}}{M_{\text{sat}} + H_{\text{app}} + H_T H_{\text{ani,oop}}}, \quad (B.19)
\]

\[
A_z = \frac{M_{\text{sat}} H_T (H_{\text{ani,oop}}^2 - H_{\text{app}}^2)}{H_{\text{ani,oop}}(M_{\text{sat}} + H_{\text{app}}) + H_T H_{\text{app}}}. \quad (B.20)
\]

The term \( H_{\text{ani,oop}} \) comes from the out-of-plane layer’s equilibrium condition A.7a, in which it was simply called \( H_{\text{ani}} \). The subscript ‘oop’ is used to
emphasize that this term, appearing in the expression for the amplitude in the in-plane layer, is in fact the anisotropy field of the out-of-plane layer. A Taylor expansions around $H_{\text{app}} = 0$ yields a small field approximation:

$$A_{z,\text{ip}} \approx \frac{B_T}{\mu_0} \left[ \frac{B_T(B_T + B_{\text{ani,oop}})}{\mu_0^2 M_{\text{sat}} B_{\text{ani,oop}}} B_{\text{app}} \right] + O(B_{\text{app}}^2)$$  \hspace{1cm} (B.21)

where $B_{\text{app}} = \mu_0 H_{\text{app}}$, $B_{\text{ani,oop}} = \mu_0 H_{\text{ani,oop}}$ and $B_T = \mu_0 H_T$ all given in T.

It can be seen that the amplitude of the precession only weakly depends on the applied field. This could have been expected since it is the laser-induced spin-current that tilts the magnetization. And the polarization of this current is, for small fields, only slightly affected by $H_{\text{app}}$.

**B.2.2 Approximation for small spin currents**

Equation B.21 is now approximated for small applied field. We ignore the second-order term in the fit:

$$A_{z,\text{ip}} \approx \frac{B_T}{\mu_0} \left[ \frac{B_T(B_T + B_{\text{ani,oop}})}{\mu_0^2 M_{\text{sat}} B_{\text{ani,oop}}} B_{\text{app}} \right]$$ \hspace{1cm} (B.22)

$$\Rightarrow A_{z,\text{ip}} \approx \frac{B_T}{\mu_0} \left[ 1 - \frac{(B_T + B_{\text{ani,oop}})}{\mu_0^2 M_{\text{sat}} B_{\text{ani,oop}}} B_{\text{app}} \right].$$ \hspace{1cm} (B.23)

The temporal, unknown field $B_T$ - due to the ultrafast spin torque - is an unwanted degree of freedom in the fitting process that we want to get rid of. The constant in front of the brackets is not important, as it does not influence the shape of B.23. Furthermore, the amplitudes are measured in arbitrary units. Consequently, we have to multiply the entire expression by some constant anyway. So, we only have to eliminate the $B_T$ inside the brackets. This will be accomplished in the following analysis by showing that $B_T + B_{\text{ani,oop}} \approx B_{\text{ani,oop}}$.

According to B.20, the z-component of the precessional amplitude in the in-plane layer is given by

$$A_z = \frac{M_{\text{sat}} H_T \sqrt{H_{\text{ani,oop}}^2 - H_{\text{app}}^2}}{H_{\text{ani,oop}}(M_{\text{sat}} + H_{\text{app}}) + H_T H_{\text{app}}}. \hspace{1cm} (B.24)$$

We introduce $\phi$ as the tilting angle of $\vec{M}$ with respect to the in-plane axis. During the derivation of B.20, we assumed that $\phi$ is small. Since $A_z = M_z =$
\[ M_{\text{sat}} \sin \phi \] we can write

\[
M_{\text{sat}} \sin \phi = \frac{M_{\text{sat}} H_T \sqrt{H_{\text{ani,oop}}^2 - H_{\text{app}}^2}}{H_{\text{ani,oop}}(M_{\text{sat}} + H_{\text{app}}) + H_T H_{\text{app}}}, \tag{B.25}
\]

\[
\sin \phi = \frac{H_T \sqrt{H_{\text{ani,oop}}^2 - H_{\text{app}}^2}}{H_{\text{ani,oop}}(M_{\text{sat}} + H_{\text{app}}) + H_T H_{\text{app}}}. \tag{B.26}
\]

Next, we rewrite B.26 in order to find an upper boundary for \( H_T \):

\[
[H_{\text{ani,oop}}(M_{\text{sat}} + H_{\text{app}}) + H_T H_{\text{app}}] \sin \phi = H_T \sqrt{H_{\text{ani,oop}}^2 - H_{\text{app}}^2}, \tag{B.27}
\]

\[
H_T \left[ \sqrt{H_{\text{ani,oop}}^2 - H_{\text{app}}^2} - H_{\text{app}} \sin \phi \right] = H_{\text{ani,oop}}(M_{\text{sat}} + H_{\text{app}}) \sin \phi, \tag{B.28}
\]

\[
H_T = \frac{H_{\text{ani,oop}}(M_{\text{sat}} + H_{\text{app}}) \sin \phi}{\sqrt{H_{\text{ani,oop}}^2 - H_{\text{app}}^2} - H_{\text{app}} \sin \phi}. \tag{B.29}
\]

Assuming that \( \sin \phi \ll 1 \) yields

\[
H_T = \frac{H_{\text{ani,oop}}(M_{\text{sat}} + H_{\text{app}}) \sin \phi}{\sqrt{H_{\text{ani,oop}}^2 - H_{\text{app}}^2}}. \tag{B.30}
\]

We expect \( H_T \) to be independent of \( H_{\text{app}} \) and \( \sin \phi \). Thus, we can calculate \( H_T \) for an arbitrary applied field. For \( H_{\text{app}} = 0 \), we assume that

\[
\phi \ll 1^\circ, \tag{B.31}
\]

\[
\Rightarrow \sin \phi \ll 0,02. \tag{B.32}
\]

Calculating B.30 for \( H_{\text{app}} = 0 \) yields an upper boundary for \( H_T \). Using the measured values - \( H_{\text{ani,oop}} = 250 \text{mT} \) and \( M_{\text{sat}} = 1,1 \cdot 10^6 \text{A/m} \) - we obtain:

\[
H_T = \frac{B_{\text{ani,oop}}(\mu_0 M_{\text{sat}} + B_{\text{app}}) \sin \phi}{\sqrt{B_{\text{ani,oop}}^2 - B_{\text{app}}^2}}, \tag{B.33}
\]

\[
H_T = \frac{H_{\text{ani,oop}}(4\pi \cdot 10^{-7} \times 1,1 \cdot 10^6 + 0) \sin \phi}{\sqrt{(250 \cdot 10^{-3})^2 - 0^2}}, \tag{B.34}
\]

\[
H_T = 5,5 H_{\text{ani,oop}} \sin \phi. \tag{B.35}
\]

Using B.32 yields

\[
H_T \ll (5,5 \times 0,02) H_{\text{ani,oop}}, \tag{B.36}
\]

\[
H_T \ll 0,11 H_{\text{ani,oop}}; \tag{B.37}
\]

\[
\Rightarrow B_T + B_{\text{ani,oop}} \approx B_{\text{ani,oop}}. \tag{B.38}
\]
Now, we can write B.23 as

\[ A_{z,ip} \approx A \left[ 1 - \frac{B_{\text{ani,oop}}}{\mu_0 M_{\text{sat}} B_{\text{ani,oop}}} B_{\text{app}} \right] = A \left[ 1 - \frac{B_{\text{app}}}{\mu_0 M_{\text{sat}}} \right], \tag{B.39} \]

in which \( A \) is some constant in A/m. Finally, we can fill in \( M_{\text{sat}} = 1,1 \cdot 10^6 \) A/m to obtain an fitting function with a fixed shape:

\[ A_{z,ip} \approx A \left[ 1 - \frac{B_{\text{app}}}{4\pi \cdot 10^{-7} \times 1,1 \cdot 10^6} \right] = A \left[ 1 - 0,72 B_{\text{app}} \right]. \tag{B.40} \]
Appendix C

Influence of the applied field angle

We will now use the macro-spin model of Figure C.1 to theoretically verify that the measured phases and amplitudes are in agreement with a spin-transfer torque-induced excitation in the top layer. Figure C.1 sketches the top layer before it is excited by the laser pulse. So, since we do not have to include a spin current in our analysis yet, the dynamics of the magnetization $\mathbf{M}$ are governed by the regular LLG equation 2.9:

$$\frac{d\mathbf{M}}{dt} = -\gamma\mu_0(\mathbf{M} \times \mathbf{H}_{\text{eff}}) + \frac{\alpha}{M}(\mathbf{M} \times \frac{d\mathbf{M}}{dt}).$$

(C.1)

Furthermore, the system is considered to be in equilibrium. Hence, we can write:

$$0 = -\gamma\mu_0(\mathbf{M} \times \mathbf{H}_{\text{eff}})$$

(C.2)

$$\Rightarrow \mathbf{M} \times \mathbf{H}_{\text{eff}} = 0,$$

(C.3)

implying that $\mathbf{M}$ and $\mathbf{H}_{\text{eff}}$ are parallel in equilibrium.

Equation C.3 is now applied to Figure C.1 in order to obtain an equilibrium condition for $\mathbf{M}$. The effective field $\mathbf{H}_{\text{eff}}$ (not drawn) is the vectorial sum of the demagnetization field $\mathbf{H}_{\text{d}} = -\mathbf{M}_z$ and the applied field $H_{\text{app}}$:

$$\mathbf{H}_{\text{eff}} = \begin{pmatrix} 0 \\ H_{\text{app}} \cos \beta \\ H_{\text{app}} \cos \beta - M_z \end{pmatrix}.$$

(C.4)

The effective field tilts the magnetization out of the film plane over an
angle $\alpha$. Thus, the magnetization is given by

$$\overrightarrow{M}_{eq} = \begin{pmatrix} 0 \\ M_{sat} \cos \alpha \\ M_{sat} \sin \alpha \end{pmatrix}. \quad (C.5)$$

Substituting C.4 and C.5 into C.3 leads to an expression for the tilting angle $\alpha$. We only give the result:

$$(B_{app} \sin \beta - \mu_0 M_{sat} \sin \alpha) \cos \alpha = B_{app} \sin \alpha \cos \beta. \quad (C.6)$$

A value for $M_{sat}$ was obtained with the field-dependent measurement in section 4.5: $\mu_0 M_{sat} = 1382$ mT. Furthermore, $B_{app} = 25$ mT. Hence, we can now calculate the equilibrium angle $\alpha$ for a given applied field angle $\beta$. It is logical that the magnetization is tilted further out-of-plane for increasing values of $\beta$. Therefore, an upper bound for $\alpha$ can be found by calculating its value for the largest angle applied, i.e. $|\beta| = 20^\circ$ (see Figure 4.9). Equation C.6 then yields

$$\alpha < 0, 2^\circ. \quad (C.7)$$

So, that means that, even for the largest value of $\beta$, the magnetization direction is not significantly influenced by the off-parallel applied field.
Appendix D

Quantification of the spin currents

In this section the method for quantifying the spin-transfer torque is presented. We use the Cu measurement of Figure 4.12 as an example. It is again shown in Figure D.1. The black dashed line shows a fit to a Beaurepaire-like exponential demagnetization curve [9], which is what we would have measured without the presence of a spin-transfer torque. The deviation from this curve is caused by a spin-transfer torque. This deviation can be obtained by the operation

$$y_{\text{STT}}(t) = y_{\text{raw}}(t) - y_{\text{decay}}(t), \quad (D.1)$$

with $y_{\text{raw}}$ the raw data points and $t$ the time in ps. The normalized result of this operation can be seen in Figure 4.13.

Next, the deviation is quantified. The distances to the lowest point under and the highest point above the black curve are denoted by respectively $y_{\text{min}}$ and $y_{\text{max}}$ (Figure D.1). We define the deviation as

$$\Delta = \left\{ \left[ \frac{y_{\text{max}} - y_{\text{decay}}(t_{\text{max}})}{y_{\text{decay}}(t_{\text{max}})} \right] + \left[ \frac{y_{\text{decay}}(t_{\text{min}}) - y_{\text{min}}}{y_{\text{decay}}(t_{\text{min}})} \right] \right\} \times 100\%, \quad (D.2)$$

in which $y_{\text{decay}}(t_{\text{max}})$ is the value for $y_{\text{decay}}$ at the time $t$ for which $y = y_{\text{max}}$ and similar for $y_{\text{decay}}(t_{\text{min}})$. So, in fact, we are expressing the deviation as a percentage of the thermal background $y_{\text{decay}}$. The resulting percentage is a measure for the amount of spin-polarized current reaching the top layer. Percentages for different spacer layer thicknesses are shown in 4.14.
Figure D.1: Quantification of the spin-transfer torque-induced contribution to the total time-resolved Kerr rotation. The dashed black line is a guide to the eye of the exponential decay contribution due to thermal dissipation.
Bibliography


