TIME- AND SPATIALLY-RESOLVED MAGNETIZATION DYNAMICS INDUCED BY SPIN-ORBIT TORQUES IN THIN FILM DEVICES

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Abstract

Conventional electronics relies on the electron’s charge as a means to transmit information. In contrast, spintronics exploits both the electron’s charge and the spin degree of freedom, enabling new functionalities such as the storage and processing of data on a single chip. The spin degree of freedom can be capitalised by lateral injection of a charge current into a heavy metal/ferromagnet bilayer system characterized by strong spin-orbit coupling. This injection geometry has proven to be an efficient spin-source, which manifests as an interfacial spin-accumulation and can be exploited as an effective magnetic field acting on the ferromagnetic layer. The resulting magnetic torques, referred to as spin-orbit torques (SOTs), allow for efficient manipulation of the magnetization on the sub-ns time scale, thereby enabling fast switching of magnetic bits in non-volatile logic units and memory applications.

In this thesis we investigate the dynamics of current-induced switching in the archetype out-of-plane magnetized Pt/Co/AlO$_x$ structures. We describe the fabrication of sub-micron Pt/Co/AlO$_x$ dots of various shapes and dimensions as well as methods to characterize and simulate their magnetic properties. We then focus on time- and spatially resolved magnetization switching measurements of circular Pt/Co/AlO$_x$ dots by scanning transmission x-ray microscopy. These measurements reveal that the magnetization reversal takes place by a fast current-induced domain wall nucleation at the edge of the Co dot with subsequent domain wall propagation of a tilted domain wall. The nucleation point is observed to alternate between the four dot quadrants, which is related to the signs of the current, magnetization, and external field. Furthermore, our measurements point out the importance of the Dzyaloshinskii-Moriya interaction and the SOTs in breaking the magnetic symmetry and thereby defining the switching polarity. The demonstration of reproducible switching for more than $10^{12}$ cycles and the negligible incubation time observed in our measurements are promising for the development of SOT based random-access memories. Our findings are supported by all-electrical switching measurements and micromagnetic simulations.

Additionally, we report on asymmetric current-induced domain wall motion in out-of-plane magnetized heavy metal/ferromagnet systems. By means of micromagnetic simulations, we find that the current-induced domain wall velocity is strongly related to the interplay of the SOTs and the Dzyaloshinskii-Moriya in-
teraction, which determines the internal non-equilibrium domain wall structure. Our results show that the domain walls are tilted and propagate asymmetrically depending on the angle between the domain wall and the current direction. Which implies that the domain wall velocity is strongly correlated with the tilt angle. As the faster domain walls are expelled rapidly from the sample boundaries, the type of domain walls experimentally observed effectively depends on the timescale of the observations. We propose a simple micromagnetic model of the asymmetric domain wall motion and find, in agreement with simulations, that the asymmetry can be controlled by applying a bias-field perpendicular to the current direction and that it increases with the injected current amplitude.
Zusammenfassung


In der vorliegenden Arbeit untersuchen wir die Dynamik von strominduziertem Magnetisierungsschalten in der Prototypstruktur Pt/Co/AlO\textsubscript{x} mit senkrechter Magnetisierung. Wir beschreiben die Herstellung von sub-Mikrometer Pt/Co/AlO\textsubscript{x} Inseln mit verschiedenen Formen und Größen und Methoden für die Charakterisierung und Simulation ihrer magnetischen Eigenschaften. Im Weiteren führen wir zeitlich und örtlich aufgelöste Experimente des Magnetisierungsschalten von runden Pt/Co/AlO\textsubscript{x} Inseln mittels Rastertransmissionsröhrenstrahlenuhrmikroskop durch. Diese Messungen zeigen, dass das Schalten durch eine schnelle Domänenwandnukleation am Rand der Co Insel mit anschließender Domänenwandpropagation einer gekippten Domänenwand erfolgt. Der Nukleationspunkt alterniert zwischen den vier Quadranten in Abhängigkeit von Strom, Magnetisierungsrichtung und externem Feld. Zusätzlich belegen unsere Messungen den Stellenwert der Dzyaloshinskii-Moriya Wechselwirkung und der SOTs. Im Zusammenspiel führen diese Beiträge zu einer Brechung der Magnetisierungssymmetrie und definieren dadurch die Schaltpolarität. Unsere Experimente belegen eine zuverlässige Magnetisierungsumkehr über mehr als 10^{12} Zyklen. Zusammen mit der Beobachtung einer vernachlässigbaren Nukleationszeit sind unsere Ergebnisse eine vielversprechende Grundlage für die Entwicklung von SOT-basierten Schreib-Lese-Speichern (engl.: random-access memory, RAM). Wir bekräftigen unsere Ergebnisse durch elektrische Messungen und mikromag-
netische Simulationen.
Acknowledgements

The doctoral studies are similar to a sailing trip with calm and stormy weather, sunny and rainy days, thrilling and boring legs. I am grateful to Prof. Dr. Pietro Gambardella who made me experience all of this and helped me to get through tough times. Not only was he a great supervisor and teacher who always had an open door for a “five minute” discussion, which normally did not end before an hour had passed, but he also was, and still is, a source of inspiration with his professionalism, worldly wisdom and perfectionism. A big thank you goes to Dr. Kevin Garello who introduced me to the field of spintronics and rock climbing, accompanied me through my master project and thesis, and took good care of me in the beginning of my Ph.D. studies. He is always available for a discussion of personal or scientific matters, which makes me very fortunate.

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Zurich, Spring 2018

M. B.
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<th>Description</th>
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<tbody>
<tr>
<td>AHE</td>
<td>Anomalous Hall effect</td>
</tr>
<tr>
<td>ANE</td>
<td>Anomalous Nernst effect</td>
</tr>
<tr>
<td>CIDWM</td>
<td>Current-induced domain wall motion</td>
</tr>
<tr>
<td>DL</td>
<td>Damping-like (torque)</td>
</tr>
<tr>
<td>DMI</td>
<td>Dzyaloshinskii-Moriya interaction</td>
</tr>
<tr>
<td>DW</td>
<td>Domain wall</td>
</tr>
<tr>
<td>EBL</td>
<td>Electron beam lithography</td>
</tr>
<tr>
<td>FL</td>
<td>Field-like (torque)</td>
</tr>
<tr>
<td>FM</td>
<td>Ferromagnet( ic)</td>
</tr>
<tr>
<td>GMR</td>
<td>Giant magnetoresistance</td>
</tr>
<tr>
<td>HM</td>
<td>Heavy metal</td>
</tr>
<tr>
<td>IPA</td>
<td>Isopropyl alcohol</td>
</tr>
<tr>
<td>LLG</td>
<td>Landau-Lifshitz-Gilbert equation</td>
</tr>
<tr>
<td>MO&lt;sub&gt;x&lt;/sub&gt;</td>
<td>Metal oxide</td>
</tr>
<tr>
<td>MOKE</td>
<td>Magneto-optical Kerr effect</td>
</tr>
<tr>
<td>MR</td>
<td>Magnetoresistance</td>
</tr>
<tr>
<td>MRAM</td>
<td>Magnetic random-access memory</td>
</tr>
<tr>
<td>MTJ</td>
<td>Magnetic tunnel junction</td>
</tr>
<tr>
<td>NM</td>
<td>Non-magnetic (metal)</td>
</tr>
<tr>
<td>OHE</td>
<td>Ordinary Hall effect</td>
</tr>
<tr>
<td>PHE</td>
<td>Planar Hall effect</td>
</tr>
<tr>
<td>PMA</td>
<td>Perpendicular magnetocrystalline anisotropy</td>
</tr>
<tr>
<td>SMD</td>
<td>Surface-mount device</td>
</tr>
<tr>
<td>SOC</td>
<td>Spin-orbit coupling</td>
</tr>
<tr>
<td>SOT</td>
<td>Spin-orbit torque</td>
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<tr>
<td>ST-FMR</td>
<td>Spin-torque ferromagnetic resonance</td>
</tr>
<tr>
<td>STT</td>
<td>Spin-transfer torque</td>
</tr>
<tr>
<td>STXM</td>
<td>Scanning transmission x-ray microscopy</td>
</tr>
<tr>
<td>TMR</td>
<td>Tunneling magnetoresistance</td>
</tr>
<tr>
<td>XMCD</td>
<td>X-ray magnetic circular dichroism</td>
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1 Introduction

The amount of digital data volume increases year by year at an astounding pace, which requires the data storage capacity to grow at least at the same speed. Even though mass storage may seem abstract to the public as cloud storage gives the impression of infinite capacity, data still has to be stored physically. Astonishingly, today’s mass storage technology is based on technology that is several decades old: hard disc drives. Their foundations have been laid out more than forty years ago with the development of advanced nano-fabrication methods for coating and thin film deposition such as molecular beam epitaxy and dc-magnetron sputtering.

These technologies have enabled the fabrication of sub-nm thick ferromagnetic (FM) films, which in turn have formed the basis of modern hard disc drives. Working with such ultrathin layers, Peter Grünberg and his colleagues discovered in 1986 that, due to the interlayer exchange interaction, two FM layers separated by a thin non-magnetic metal (NM) spacer align parallel (ferromagnetic) or antiparallel (antiferromagnetic) depending on the thickness of the spacer layer. Measurements of the electrical resistance along the sample plane, as well as perpendicular to the sample plane of such a trilayer system revealed a giant change in the magnetoresistance, as a function of the relative orientation of the two FM layers (Fig. 1.1a and b). The discovery of this giant magnetoresistance (GMR) effect in 1988 by A. Fert and P. Grünberg, for which they received in 2007 the Nobel prize in physics, marks the birth of spintronics. The GMR effect results from spin dependent scattering in the two FM layers (Fig. 1.1c and d). Such spin-valve structures – the name relates to the physical origin of the GMR effect – are still used today as magnetic field sensors due to their high sensitivity to magnetic fields, e.g. in read heads of hard disc drives.

In 1996, Slonczewski and Berger predicted spin-transfer torque (STT) as a promising route to manipulate the magnetization of spin-valves. They anticipated that a spin polarized current will exert a torque on the magnetic moments of a ferromagnet through which it is flowing (see Fig. 1.2a) by means of s-d exchange interaction. One of the most important implications of their theory is that a current flowing perpendicularly through a spin-valve structure will get spin-polarized by the first “fixed” ferromagnetic layer (FM1) and may exert a torque on the second “free” ferromagnetic layer (FM2) Ideally, STT is sufficient to switch the magnetization of FM2 by transferring spin angular mo-
Figure 1.1 | Giant magnetoresistance in spin-valves. a, Interlayer exchange coupling in Fe/Cr/Fe trilayers showing antiferromagnetic coupling of two Fe-layers separated by a Cr-spacer and b, the corresponding electrical measurements revealing a GMR effect of 1.5% (Figures adapted from Grünberg et al.\(^5\)). c, Schematic representation of a spin-valve structure in the perpendicular current geometry. d, Illustration of the GMR effect based on spin dependent scattering and the two current model equivalent circuit. Ferromagnetic (antiferromagnetic) alignment of FM1 and FM2 results in a low (high) resistance state.

\(\delta R(R_\text{e})\)
This discovery paved the way for the magnetic racetrack memories proposed by Parkin et al.\textsuperscript{19}. In such memories, digital data is encoded by means of magnetic domains which can be pushed through the racetrack by STT.

In parallel to the spintronics development boosted by the GMR and STT, the discovery of the tunnelling magnetoresistance (TMR) made by Julliere et al.\textsuperscript{20} and Tedrow et al.\textsuperscript{21} in the 70’s did not attract significant attention until the first successful measurements at room temperature were published in 1995\textsuperscript{22,23}. In contrast to the GMR, which is observed in trilayers consisting of a NM sandwiched between two FM layers, TMR is observed in trilayers consisting of an insulating metal oxide (MO\textsubscript{x}) layer sandwiched between two FM layers. Such FM/MO\textsubscript{x}/FM heterostructures are referred to as magnetic tunnel junctions.
(MTJs). Similar to the spin-valves, a current flowing through a MTJ (perpendicular to the sample plane) experiences a different resistance, depending on the relative orientation of the two FM s (Fig. 1.3c). However, the underlying physics is very different from the spin dependent scattering causing the GMR. The origin of the TMR is spin-conserving tunneling through the MO \textsubscript{x} layer, where the tunneling probability, and therefore the electrical resistance, depends on the relative orientation of the two FM layers. Hence, in contrast to GMR, TMR is a purely quantum mechanical effect. The relative resistance change of a MTJ (the TMR) is defined as $TMR = (R_{AP} - R_P)/R_P$, with $R_P$ ($R_{AP}$) being the resistance of the parallel (antiparallel) magnetization state. For over 20 years little interest has been shown by the spintronics community in Jullieres measurements. This is partly due to the technologically very demanding fabrication process for the production of reliable MTJs, as well as due to the relatively small effect, which could only be measured at 4 K at the time.

The development of MTJs was intensified after the prediction\cite{24,25} and measurements\cite{26} of large TMR ratios in MgO-based MTJs and finally by the demonstration of more than 600\% TMR at room temperature in CoFeB/MgO/CoFeB tunnel junctions in 2008\cite{27}. Since the mid 2000s, state of the art read heads in hard disc drives are based on MTJs\cite{28} and completely replaced GMR read heads. Moreover, these promising results and the significantly increased read-out signal compared to all-metallic spin-valves led to the development of MTJ-based STT-MRAM\cite{16}. However, the development of STT-MRAMs is hampered by the very complex fabrication process as ultrathin oxide-barriers are required for efficient writing. Alternative spin-current sources, which would no longer require a spin-polarizer and thereby reduce the requirements on the oxide-barriers, were well perceived by the spintronics community.

Two different effects, both based on spin-orbit coupling, were identified as promising spin-sources; the spin-Hall effect (SHE), theoretically predicted in 1971 by Dyakonov and Pere\cite{31}, where spin dependent scattering of conduction electrons in materials with a large spin-orbit coupling (mainly heavy metals) leads to a spin current orthogonal to the electron flow, and the relativistic Rashba-Edelstein\cite{32,33} effect, which induces a spin-polarization in systems with broken inversion symmetry. Underlining the link to their physical origin, torques arising from these effects were named spin-orbit torques (SOTs). In 2009 Chernyshov et al.\cite{34} showed domain rotation and hysteretic switching by means of SOTs in ferromagnetic semiconductors. Shortly thereafter, Miron et al. showed sizeable SOTs in Pt/Co/AlO\textsubscript{x} trilayers with perpendicular magne-
Figure 1.3 | SOT-induced magnetization switching. a, The lateral injection of a current into a HM/FM/MO$_x$ trilayer enables to switch the FM between its two remanent states by means of SOTs. b, First report of SOT-induced magnetization switching by Miron et al.\textsuperscript{29} in 2011 on a Pt/Co/AlO$_x$ trilayer. c, Schematic illustration of an out-of-plane magnetized MTJ for SOT-induced switching. The injection of a current pulse through the current line (red) allows to control the state of the free layer (purple). The inset shows schematically the TMR as a function of out-of-plane field and the state of the two FM layers. d, First experimental demonstration of SOT-induced writing of a perpendicular MTJ and read-out by the TMR. (Figures adapted from Refs. 29 and 30)

tocrystalline anisotropy, which enabled current-induced domain wall motion\textsuperscript{35,36} and eventually magnetization switching\textsuperscript{29}. These experiments mark the evolution of \textit{spintronics} towards \textit{spinorbitronics}, where spin currents are generated by means of spin-orbit interaction. Hence, spinorbitronics does no longer require technically demanding spin-valves or MTJs to generate spin-currents, as easy to fabricate heavy metal (HM)/FM-heterostructures can be used as efficient spin-current sources (see Figs. 1.3a and b).

These first demonstrations of SOT-induced magnetization dynamics triggered significant interest in the spintronics field and led to the investigation of spin-orbit coupling phenomena in numerous HM-based material systems (Ta/-CoFeB/MgO, Pt/Co/AlO$_x$, Ti/CoFeB/Pt, W/CoFeB/Ti, etc.), in antiferromagnetic systems (IrMn/NiFe/TaN, PtMn/Co/MgO, CuMnAs, etc.) and in more
exotic materials such as Heusler alloys and topological insulators.\cite{37,47} In 2014, Garello et al.\cite{48} demonstrated sub-ns switching of 100 nm large Co dots, which triggered large interest in the MRAM community. Furthermore, the demonstration that the bit-state of a Ta/CoFeB/MgO/CoFeB MTJ can be written by means of SOT, using an in-plane current injection geometry,\cite{30} proved the feasibility of the proposed SOT-MRAM\cite{49} (see Fig. 1.3c and d). The in-plane current injection geometry, inherent to SOT based devices, has the advantage of decoupling the read and write lines, thereby reducing electrical stresses through the tunnel barrier. However, such three terminal devices have the disadvantage of a decreased packaging density due to the increased number of required electrical connections. In parallel to the SOT-MRAM development, SOT-induced domain wall motion has been investigated extensively.\cite{36,50–53} This pinpointed the importance of the Dzyaloshinskii-Moriya interaction,\cite{54–57} an additional exchange term, in stabilizing chiral Néel domain walls characterized by very high mobilities. Recently, in antiferromagnetically coupled stripes current-induced domain wall speeds of up to 750 m s$^{-1}$ have been observed,\cite{58} which is expected to further enhance the development of racetrack memories.\cite{59}

Today, the underlying effects leading to SOTs, as well as their manifold manifestations and related effects are strongly investigated all over the world through diverse approaches. Significant research was focused on the SOT-induced magnetization switching processes, both experimentally,\cite{48,60–62} and theoretically.\cite{63–66} Nevertheless, a number of issues remain outstanding and require further investigation. Notable open questions concern the timescale and physical processes of the magnetization switching during the current injection, the role played by the SOTs and the Dzyaloshinskii-Moriya interaction, as well as the intermediate magnetization configurations. The present work addresses these remaining issues and sheds light on the micromagnetic reversal process during SOT-induced magnetization switching in the archetype out-of-plane magnetized heterostructure Pt/Co/AlO$_x$ by means of time- and spatially resolved x-ray measurements. The deeper understanding of the reversal process is not only interesting from a fundamental point of view, but may also pave the way for improved SOT-MRAM functionality.

This thesis is organized as follows. In Chapter 2 we will introduce the important magnetic interactions, as well as the spin-orbit torques and related phenomena. In Chapter 3 the main methods used in this thesis are presented: device fabrication, electrical measurements and analysis, x-ray based measurements,
and micromagnetic simulations. The focus of Chapter 4 is on time- and spatially
resolved magnetization dynamics by means of electrical and x-ray based mea-
surements. The presented findings are underlined by micromagnetic simulations.
In Chapter 5, we evidence asymmetric current-induced domain wall motion in
HM/FM bilayers with out-of-plane magnetization. Finally, we will conclude and
give an outlook on possible future studies in Chapter 6.
2 Background

2.1 Basic magnetic interactions

In the following section we introduce the relevant magneto- and electrostatic interactions acting on the magnetization \( M \) of a FM. We will see, that compared to the long ranged and rather weak magnetostatic interactions (e.g. the interaction of a magnet with an external magnetic field) the short-ranged electrostatic effects, like the exchange interaction, can be very strong and dominant.

2.1.1 Zeeman energy

The interaction of a magnetic material with a magnetic field is described by the Zeeman energy \( E_Z \). It is the potential energy of a magnetic volume \( V \) exposed to an external field \( H_{\text{ext}} \) and is given by

\[
E_Z = -\mu_0 \int_V M \cdot H_{\text{ext}} \, dV, \tag{2.1}
\]

where \( \mu_0 = 4\pi 10^{-7} \text{ J A}^{-2}\text{m}^{-1} \) is the vacuum permeability. For a single magnetic moment \( m \), this translates to the microscopic energy

\[
E_Z = -m \cdot B_{\text{ext}}, \tag{2.2}
\]

with \( B_{\text{ext}} \) being the local magnetic field, defined as \( B_{\text{ext}} = \mu_0 H_{\text{ext}} \). Obviously, \( E_Z \) is minimized if \( m \) and \( B_{\text{ext}} \) align parallel. This can be described by the magnetic torque \( T_Z = m \times B_{\text{ext}} \) acting on \( m \), as we illustrate in Fig. 2.1. \( T_Z \) will induce the magnetic dipole to rotate and align it with \( B_{\text{ext}} \), indicated by the two small arrows on top and bottom of the dipole. Any magnetic torque \( T \) can equivalently be described by an effective magnetic field \( B_T = T \times m \) which rotates with \( m \).
2.1. Basic magnetic interactions

Figure 2.1 | Magnetic interactions. a, Illustration of a the Zeeman torque $\mathbf{T}$ acting on a magnetic dipole $\mathbf{m}$ exposed to a constant magnetic field $\mathbf{B}$. The magnetic dipole tends to align with the magnetic field. b, Schematic representation of the shape anisotropy effect caused by the demagnetizing field of an elongated thin film. Energetically favoured is the magnetization configuration for which $\mathbf{H}_{\text{dem}}$ (and the free magnetic charges) is minimized; in-plane magnetization shown in the upper panel. The out-of-plane configuration (lower panel) is not favoured by the shape anisotropy. c, Illustration of the ferromagnetic (antiferromagnetic) exchange interaction for $J > 0$ ($J < 0$) according to Eq. (2.3). d, Schematic of the interfacial DMI caused by the large SOC of the HM layer (green). The DMI tries to maximize the angle between two adjacent spins of the FM against the exchange interaction and the uniaxial anisotropy energy. $d > 0$ ($d < 0$) stabilizes left (right)-handed spin spirals and DWs.

2.1.2 Exchange interaction

The short ranged exchange interaction between two atomic spins arising from the Coulomb interaction and the Pauli principle is often described by the Hamil-

\[ \text{Actually } \mathbf{B}_{\text{ext}} \text{ is the magnetic induction defined as } \mathbf{B}_{\text{ext}} = \mu_0 (\mathbf{M} + \mathbf{H}_{\text{ext}}) \text{ and } \mathbf{H}_{\text{ext}} \text{ the magnetic field generated e.g. by an electromagnet. In our experimental setup we regard the space between the poles of the electromagnet as vacuum (} \mathbf{M} = 0 \text{ and hence } \mathbf{B}_{\text{ext}} = \mu_0 \mathbf{H}_{\text{ext}} \text{ which for simplicity we call magnetic field.} \]
2.1. Basic magnetic interactions

Hamiltonian

\[ \mathcal{H}_{ex} = -2J \mathbf{S}_1 \cdot \mathbf{S}_2. \]  \hspace{1cm} (2.3)

Herein \( J \) is the exchange constant with dimension of an energy and \( \mathbf{S}_{1,2} \) are the dimensionless unit vectors representing the two atomic spins. As we illustrate in Fig. 2.1b, \( J > 0 \) (\( J < 0 \)) indicates a ferromagnetic (antiferromagnetic) interaction which tends to align the spins parallel (antiparallel). If we consider a lattice of spins (e.g. a crystal) the Hamiltonian is generalized to

\[ \mathcal{H}_{ex} = - \sum_{i \neq j} J_{ij} \mathbf{S}_i \cdot \mathbf{S}_j. \]  \hspace{1cm} (2.4)

In this Hamiltonian \( \mathbf{S}_{i,j} \) represents the atomic spin on the lattice sites \( i, j \). This Hamiltonian can be further simplified if only nearest neighbour interactions are considered. The simplicity of Eq. (2.4) is the basis of its success and the reason why it is still widely used today despite its limitations to describe systems with delocalized electrons such as Co, Fe and Ni.

2.1.3 Demagnetizing field

It is well known that two (or more) separated magnetic bodies interact with each other (e.g. earth and compass). This interaction has its origin in the magnetic field created by the magnetization of these objects. The surface magnetization of a ferromagnet can mathematically be described by free “magnetic charges” with a density \( \rho_m = -\nabla \cdot \mathbf{M} \) due to the discontinuity of \( \mathbf{M} \) at the boundary between two materials (vacuum and the magnet itself). In the upper panel of Fig. 2.1b this is illustrated for a elongated thin film where the magnetization lies in the sample plane. The magnetic charges are traditionally associated with the north (+) and south (-) pole. Following Ampère-Maxwell’s law, we have \( \nabla \cdot \mathbf{H} = -\nabla \cdot \mathbf{M} \). Hence, the magnetic charges induce a magnetic field which exists outside and inside of the magnet. In the former case this field is called stray or dipolar field \( \mathbf{H}_{dip} \) and is the one responsible for the alignment of a compass in the earth’s magnetic field. In the latter case this field is called demagnetizing field \( \mathbf{H}_{dem} \) as it opposes the internal magnetization. Both fields, \( \mathbf{H}_{dip} \) and \( \mathbf{H}_{dem} \), depend on \( \mathbf{M}(\mathbf{r}) \) and the shape of the sample as this defines the magnetization discontinuity and therefore the distribution of the free magnetic charges. For an uniformly magnetized sample with an elliptical shape, \( \mathbf{H}_{dem} \) is uniform and can
be calculated analytically

\[ H_{\text{dem},i} = -N_{ij}M_j \quad i, j = x, y, z, \tag{2.5} \]

where \( N_{ij} \) is the demagnetizing tensor with trace \( N_{xx} + N_{yy} + N_{zz} = 1 \). For some simple shapes (e.g. long needle, sphere, thin film, etc.) \( H_{\text{dem}} \) can be approximated by symmetry arguments starting from Eq. (2.5). However, for more complicated shapes \( H_{\text{dem}} \) cannot be calculated analytically anymore and has to be approximated numerically.

From an application point of view, it is important to note that \( H_{\text{dem}} \) induces a magnetic anisotropy. According to Eq. (2.1), the directions of \( M \) are favoured which minimize \( H_{\text{dem}} \). As \( H_{\text{dem}} \) depends on the sample shape, this effect is called shape anisotropy. For the thin films investigated herein, shape anisotropy favours the magnetization directions in the sample plane. The out-of-plane magnetization configuration illustrated in the lower panel of Fig. 2.1b is energetically unfavoured due to the increased \( H_{\text{dem}} \) originating from the large amount of magnetic surface charges.

On top of the shape anisotropy, \( H_{\text{dem}} \) is responsible for multi domain magnetization configurations such as magnetic flux closure states, which can further reduce the energy of a system (at the cost of an increase in \( E_{\text{ex}} \)) by balancing all the magnetic charges.

### 2.1.4 Magnetocrystalline anisotropy energy

The magnetization of a ferromagnetic domain can be constrained to lie preferentially along specific crystal directions called easy axes. Thus, the associated magnetocrystalline anisotropy energy depends on the crystal structure. For hexagonal and tetragonal crystals (e.g. Co) the leading term of the magnetocrystalline anisotropy energy \( E_a \) is

\[ E_a = K_2 \sin^2(\theta), \tag{2.6} \]

in units of J m\(^{-3}\). \( K_2 \) is the anisotropy constant and \( \theta \) the angle between the magnetization direction \( M \) and the easy axis. Depending on the crystal symmetry, additional higher order terms have to be considered (i.e. \( K_4 \sin^4(\theta) \)). The present work focuses on samples exhibiting an uniaxial out-of-plane easy axis along \( \pm z \) referred to as perpendicular magnetocrystalline anisotropy (PMA). In ultrathin FM layers the interfacial out-of-plane anisotropy can exceed the
in-plane volume anisotropy, and favour out-of-plane magnetization\footnote{\cite{footnote}}. Such an anisotropy field can be parametrized as an effective magnetic field $B_k = B_k (M \cdot z)$, $z$ which points along $\pm z$ depending on the sign of $M_z$. The shape anisotropy opposes the PMA. Therefore, the measured effective anisotropy field $B_k$ used to model the uniaxial out-of-plane anisotropy is $B_k = B_a - B_{dem}$, with $B_a$ being the “real” anisotropy field given by Eq. (2.6) and $B_{dem} = \mu_0 H_{dem}$.

### 2.1.5 Dzyaloshinskii-Moriya interaction

In contrast to the exchange interaction described above, the interfacial Dzyaloshinskii-Moriya interaction (DMI) is an antisymmetric exchange interaction, which tries to couple two adjacent spins perpendicularly. The DMI-Hamiltonian is given by

$$H_{DM} = -D_{ij} \cdot (S_i \times S_j), \quad (2.7)$$

with $D_{ij} = d (u_{ij} \times z)$ being the DMI vector and $S_i,j$ two neighbouring atomic spins connected by the vector $u_{ij}$ (see Fig. 2.1d). Due to the interfacial origin of the DMI, $d$ has the dimension $J \, m^{-2}$ and was measured to be of the order of $1 \, \text{mJ} \, \text{m}^{-2}$ (Ref. \cite{footnote}). If we neglect anisotropies and external magnetic fields, we see that the actual angle between two neighbouring spins results from the competition of $H_{ex}$ and $H_{DM}$. If we compute the spin configuration for a chain of localized spins, as we do in the upper panel of Fig. 2.1d for $d > 0$, we notice that $H_{DM}$ is a chiral interaction which stabilizes a left-handed spin spiral when following the path of the spin chain. In line with this argument, $d < 0$ stabilizes right-handed spin spirals, shown in the lower panel of Fig. 2.1d. With the same argumentation one can predict that the DMI will stabilize chiral Néel domain walls (see next section) and topologically protected quasi-particles such as skyrmions\cite{footnote}\cite{footnote}.

### 2.1.6 Domains and domain walls

In extended films magnetic domains form due to the demagnetizing field. The size of the domains is determined by the competition between the energy associated with the domain wall (DW) and the demagnetizing field. The DW width $\delta_{DW}$ is determined by the competing exchange interaction and the effective anisotropy\footnote{\cite{footnote}}. $\delta_{DW}$ is maximized by the exchange interaction, which tends to

\footnote{The interfacial contribution to the anisotropy energy\cite{footnote} has it origin in the so-called spin-orbit coupling, which will be introduced later. On the other hand, the volume anisotropy contribution originates mainly from the demagnetizing field introduced in the previous section.}
align neighbouring moments parallel. On the other hand, the effective anisotropy energy will be minimized if all the magnetic moments are pointing along the easy axes of the system, which favours small $\delta_{DW}$. The wall width can be approximated as $\delta_{DW} \approx \sqrt{A/K_{eff}}$ with the effective anisotropy constant $K_{eff}$, which takes the demagnetizing field into account. $A$ is the exchange stiffness related to the exchange constant $J$ (Ref. [52, 75]). The corresponding DW energy ($E_{DW} \approx \sqrt{AK_{eff}}$) amounts typically to 1 mJ m$^{-2}$, which is the same order of magnitude as the energy associated with the DMI. In Fig. 2.2 we illustrate two common DW types; the Bloch wall and the Néel wall.

**Bloch wall**

In the upper (lower) panel of Fig. 2.2a we illustrate a right (left)-handed Bloch wall, in which the magnetization spirals in the plane of the DW as we cross it. In wide FM thin films characterized by strong PMA, Bloch walls are the lowest energy DWs (see Fig. 2.3a) as they do not induce divergence of the magnetization, $\nabla \cdot \mathbf{M} = 0$, and hence minimize magnetic volume charges. However, in narrow thin films, a Bloch wall creates additional surfaces charges and therefore its energy increases.
2.1. Basic magnetic interactions

Figure 2.3 | Dependence of the domain wall configuration. a, Phase diagram of the DW configuration as a function of the thickness and the width of a ferromagnetic stripe with PMA. The DMI is not taken into account. b, Phase diagram of the azimuthal domain wall angle $\psi$ as a function of the wire width. A pure Néel (Bloch) wall corresponds to $\psi = 0^{\circ}$ ($90^{\circ}$). The green line shows that the DMI stabilizes intermediate DW configurations. (Figures adapted from Refs. 76 and 77).

Néel wall

In contrast to the Bloch wall, a Néel wall is characterized by a magnetization rotation perpendicular to the DW (see Fig. 2.2b). In a bulk FM a Néel wall has higher energy than a Bloch wall as $\nabla \cdot \mathbf{M} \neq 0$. However, as we approach the limit of thin and narrow wires, Néel walls are energetically favoured by magnetostatics as they do not create any surface charges and have no associated $\mathbf{H}_{\text{dip}}$. The DMI can further lower the energy of Néel walls and increase their stability in ultrathin films. Contrary to the general consensus, it was recently shown that, in perpendicularly magnetized wires, intermediate DW states, neither pure Néel nor Bloch walls, can be stabilized even in the absence of DMI, if the dimensions are chosen properly (see Fig. 2.3b).

The energy of a Néel wall is typically independent of its sense of rotation (left- or right-handed). In HM/FM thin films, however, the degeneracy of left- and right-handed DWs is lifted by the DMI as we mentioned in Section 2.1.5. In out-of-plane magnetized Pt/Co/AlO$_x$ layers left-handed Néel DW were found to be stabilized by the DMI.

2.1.7 Spin-orbit torques

As mentioned in the introductory chapter, recently, a writing scheme based on an in-plane current injection geometry, exploiting the spin-orbit coupling in heavy
metals, was shown to be an efficient way to control the magnetization orientation\cite{29,37,48,79}. The following section focus on this new mechanism referred to as spin-orbit torques.

**Origin**

When a current is injected laterally into a trilayer consisting of a heavy metal (HM)/ferromagnet (FM)/metal oxide (MO$_x$) the charge current at the HM/FM interface becomes spin polarized. The origin of the spin polarization is believed to be the spin-Hall effect\cite{31} (SHE) or the relativistic Rashba-Edelstein\cite{32,33,80} effect. In both cases the underlying effect is based on the coupling between the electron spin and the orbital motion, the spin-orbit coupling\cite{81} (SOC), which is large in a HM.

The spin-Hall effect (intrinsic and extrinsic) arises from spin-dependent scattering in the bulk of the HM, which deviates electrons with different polarization in different directions. This results in a spin-current perpendicular to the charge current\cite{82,83} (see Fig. 2.4a), which in turn induces a non-equilibrium spin accumulation at the HM/FM interface. The spin polarization builds up (or decays) exponentially, which is characterized by the material dependent spin diffusion length $\lambda_s$ (Ref. 84). $\lambda_s$ is several hundreds of nm in materials with weak SOC (semiconductors\cite{85} and some light metals like Cu\cite{86} and Al\cite{87}). Materials characterized by a large SOC (e.g. Pt\cite{88,89}, W\cite{40} or Ta\cite{86}) show typically very short

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**Figure 2.4 | Origin of the the spin-orbit torques.**  

_a_, Spin-dependent scattering in the bulk of the HM (due to the spin-Hall effect) leads to a spin-current perpendicular to the charge current and introduces a non-equilibrium spin-accumulation at the HM/FM interface which manifests itself as a DL (and FL) SOT.  

_b_, Due to the relativistic Rashba-Edelstein effect electrons propagating at the inversion asymmetric HM/FM interface experience an effective magnetic field which can induce a net spin polarization related to the FL (and DL) SOT.
\[ \lambda_s (\sim 1 - 10 \text{ nm}) \] and are therefore the perfect candidates as spin sources for the ultrathin HM/FM bilayers investigated herein, because the comparable \( \lambda_s \) and HM-thickness leads to a large spin accumulation at the HM/FM interface. The accumulated spins are absorbed in the FM giving rise to a local torque acting on the ferromagnet.

The Rashba-Edelstein effect originates from the symmetry breaking of the HM/FM interface, which creates a net electric field perpendicular to the interface. According to the Lorentz transformation, electrons propagating at relativistic speed in such an electric field experience an effective magnetic field perpendicular to the charge current and the electric field. This effective magnetic field induces a non-equilibrium spin accumulation at the interface (see Fig. 2.4b). Similar to the SHE, the spin accumulation results in a local torque. Due to the very short screening length in metals (a couple of Å), ultrathin FM layers are important for the Rashba-Edelstein effect to manifest itself as a spin source.

Independent of the origin of the spin accumulation at the HM/FM interfaces, the resulting torques are called spin-orbit torques (SOTs), which underlines their common origin, the SOC. Consistent with the opposite sign of the SOC in Ta and Pt, the SOTs are reversed in Pt/Co/AlO\(_x\) and Ta/CoFeB/MgO.

**Symmetries**

Electrical measurements on HM/FM/MO\(_x\) thin films revealed that SOTs with distinct symmetries are acting on the magnetization (\( \mathbf{m} \)), the damping-like (DL) torque
\[ T^{DL} = T^{DL} \mathbf{m} \times (\mathbf{y} \times \mathbf{m}), \tag{2.8} \]
and the field-like (FL) torque
\[ T^{FL} = T^{FL} \mathbf{m} \times \mathbf{y}. \tag{2.9} \]

Here, \( \mathbf{y} \) is the direction orthogonal to the current direction (\( \mathbf{x} \)) and the sample normal (\( \mathbf{z} \)) as indicated in Fig. 2.4a and b. The amplitudes of the SOTs are parametrized by the factors \( T^{DL} \) and \( T^{FL} \). Using the sign convention of Eqs. (2.8) and (2.9), \( T^{DL} > 0 \) and \( T^{FL} > 0 \) (opposed to \( B_{Oe} \)) for Pt/Co/AlO\(_x\) samples. Measurements\[37\] and symmetry arguments\[93\] revealed that additional higher order terms can exist, which make the factors \( T^{DL} \) and \( T^{FL} \) depend on the angle of \( \mathbf{m} \) (see Eqs. (A.3) and (A.4)). Alternatively, the SOTs can be rewritten in terms of effective current-induced magnetic fields \( \mathbf{B}^{DL(FL)} \) using the definition
of a magnetic torque $\mathbf{T} = \mathbf{m} \times \mathbf{B}$ introduced in Section 2.1, leading to

$$\mathbf{B}^{DL} = \mathbf{B}^{DL}(\mathbf{y} \times \mathbf{m}),$$

(2.10)

and

$$\mathbf{B}^{FL} = \mathbf{B}^{FL}\mathbf{y},$$

(2.11)

with $\mathbf{B}^{DL(FL)} = T^{DL(FL)}$ and $\mathbf{m}$ the unit vector of magnetization. Due to this equivalence both, SOTs ($T^{DL(FL)}$) and effective current-induced magnetic fields ($\mathbf{B}^{DL(FL)}$), are used herein depending on the context. Figure 2.4 illustrates the effective SOT fields acting on the magnetization during the injection of a positive current.

At this point it is important to note that the SHE and Rashba-Edelstein effect can in principle induce both torques with the observed DL and FL symmetry. This led to a controversial discussion in the SOT community with the result that the DL torque is mainly attributed to the SHE and the FL torque to the Rashba-Edelstein effect. However, the separation of both effects is difficult as both depend on the HM layer thickness and spin-dependent scattering at the HM/FM interface may contribute to the spin accumulation.

**Measurement principles**

Over the last years various SOT measurement techniques were developed. All of them have in common that they measure the current-induced effective fields and use symmetry arguments to relate them to the amplitudes of the DL and FL torque.

The first method is based on SOT-induced magnetization switching measured as a function of external field. The critical switching current can be translated into an effective magnetic field by comparison with the external field $B_{ext}$ and the effective anisotropy field $B_k$. However, as relatively large currents are needed for the switching, it is assisted by Joule heating and does not show the expected monodomain behaviour required for the SOT quantification. Nevertheless, this method is the final proof for SOTs with finite amplitudes.

The second method is spin-torque ferromagnetic resonance (ST-FMR). The injection of an rf-current into the HM/FM/MO$_x$ at resonance frequency for an externally applied magnetic field causes the magnetization to precess. This leads to a periodic modulation of the magnetoresistance (MR) and hence to a rectified voltage originating from the mixing of the rf-current and the oscillating MR. Lineshape analysis of the rectified voltage as a function of the applied mag-
netic fields enables the quantification of any additional effective field such as the current-induced SOTs. Such ST-FMR measurements are typically restricted to thin films of several nm thickness.

Another method exploits the magneto-optical Kerr effect (MOKE) to measure the deviation of the magnetization from its equilibrium position due to the SOTs. A polarized laser beam with normal incident on the sample changes its polarization angle depending on the out-of-plane magnetization component. The rotation of the polarization can hence be related to the effective magnetic field acting on the magnetization and enables the quantification of the SOT amplitudes. However, in order to obtain quantitative values, the measured signals are typically calibrated with the Oersted field. Therefore, the obtained SOT amplitudes strongly depend on the Oersted field model employed.

The experimental technique of choice for this work is based on harmonic Hall voltage measurements described in detail in Section 3.3. The fundamental principle is similar to the experimental method introduced previously where the deviation of the magnetization from the equilibrium position is detected by means of the MOKE. Instead of the MOKE, harmonic Hall measurements exploit the anomalous and planar Hall effect as a very sensitive tool to track the magnetization position.

Injection of a charge current into a NM under an applied magnetic field generates the well known ordinary Hall voltage perpendicular to both the current and field direction. The ordinary Hall effect (OHE) has its origin in the Lorentz force that acts on the charge carriers and deviates them orthogonally to the current direction and the applied magnetic field. This leads to a charge separation and the creation of a transverse electric field which establishes a measurable steady state electric potential. Additionally to the OHE, in FM materials the SOC induces a transverse voltage that is proportional to the out-of-plane magnetization component \( m_z \) rather than to the magnetic field. This effect is referred to as the anomalous Hall effect (AHE). Similar to the spin-Hall effect discussed in Section 2.1.7, the SOC leads to spin dependent scattering in the FM which in turn generates an accumulation of charge carriers with distinct spin polarization at the sample edge transverse to the current direction. Due to the imbalance in the number of the electrons with up (majority) and down (minority) spins in a FM material, a transversal potential arises at the edge of the sample which depends on the \( m_z \) component of the magnetization. Lastly, the planar Hall effect (PHE) is the transverse manifestation of the anisotropic longitudinal magnetoresistance. Depending on the magnetization direction, the charge...
distribution in the atomic orbitals change. Hence the scattering cross section and the resistance experienced by the charge carriers in a FM is related to the relative orientation of the magnetization and the injected current direction. The s-d scattering thereby induces a charge accumulation transversal to the current direction. The resulting electric potential is proportional to the product of the in-plane magnetization components ($m_xm_y$).

**Spin-orbit torque efficiency**

For material systems with nominally identical composition, different SOT values have been reported. Exemplary, for the Pt(3 nm)/Co(0.6 nm)/AlO$_x$ system with PMA values for $T^{DL}$ varying between 1.7 and 7 mT, and for $T^{FL}$ between 0 and 4 mT per 10$^7$ A cm$^{-2}$, were reported. This large span of amplitudes underlines the sensitivity to the material parameters such as the magnetization, the PMA or the oxidation state and the challenges to characterize the SOTs correctly. In order to compare the SOTs between different material systems, it is useful to define the spin-torque efficiencies:

$$\xi^j_{DL(FL)} = \frac{2e}{\hbar} M_s t_{FM} \frac{B_{DL(FL)}}{j_c}, \quad (2.12)$$

where $M_s$ is the saturation magnetization, $t_{FM}$ the thickness of the FM layer and $j_c$ the current density at which the effective fields $B_{DL(FL)}$ are measured. This definition reflects that the measured effective fields scale with the volume of the magnetization and the injected current density. In order to avoid artefacts from thickness inhomogeneities, the spin-torque efficiency can alternatively be expressed per unit electric field $E = j_c\rho$ driving the injected current

$$\xi^E_{DL(FL)} = \frac{2e}{\hbar} M_s t_{FM} \frac{B_{DL(FL)}}{E}. \quad (2.13)$$

This is beneficial for voltage-controlled measurements. $\xi^E_{DL(FL)}$ relates the spin current absorbed in the FM with the driving electric field and can therefore be understood as an effective spin conductivity.

**Manifestation**

Since the first reports of sizeable SOTs in HM/FM/MO$_x$ trilayers, multiple SOT manifestations have been observed. Among the first observation was SOT-induced DW displacement in HM/FM stripes with PMA. It was also demonstrated that magnetic nano-oscillators can be driven by means of SOTs.
2.2. The art of magnetic writing

The efficient control (in terms of energy, speed, density, retention and endurance) of the magnetic bit state is key for a reliable and persistent MRAM. In the last 60 years, diverse approaches have been explored for writing magnetic bits. The most promising candidates are magnetic field, STT, and SOT. In this section we will introduce the control by magnetic field and STT. The following Section 2.3 gives an extended overview of magnetization switching by SOTs, the main focus of this thesis.

2.2.1 Writing by an external magnetic field

The precursor of modern magnetic random access memories (MRAMs) were developed in the 1950s and used for many years; the toroidal memory, also known as magnetic core memories, consisting of tiny magnetic rings (toroids) arranged around the crossing points created by a lattice of orthogonal write wires. Additionally, a single read wire passes through all the toroids in a sequential way. In such memories the digital information is stored in the magnetization sense of rotation of the toroids, being either clockwise (“1”) or anticlockwise (“0”). The state of the toroid is written by exploiting the Oersted field due to a current passing through the write wires. Depending on the combination of the wires through which the current is injected, the bit to be written is selected. The state of the magnetic bit is read by detecting an induced voltage pulse in the read wire, resulting from a change in the magnetization sense of rotation. For this, a write pulse of known polarity is injected (e.g. writing a “1”) and the resulting voltage pulse in the read wire is detected. If no voltage pulse is detected, the bit state was “1” already, if a voltage pulse is detected the bit was a “0”. Hence, the bit state may be changed during the read process and has to be updated after each access.

Recent MRAM technology manages to access data without such a destructive III From Biot-Savart law it is known that a current passing through a wire generates a circular magnetic field around it, known as Oersted field.
read process due to the development of spin-valves and magnetic tunnel junctions (see introductory chapter). In these MRAMs the readout is done by measuring the resistance of the MTJ, which is different (several 100 %) for the parallel, low resistance state ("1") and the antiparallel, high resistance state ("0") thanks to the TMR\textsuperscript{27}. However, the main challenge remains: how to efficiently control the relative orientation of the two FM layers of an MTJ. The obvious solution is to control it by means of an external field, which led to the development of field-MRAM\textsuperscript{122,123}.

Field-MRAM are comparable to the core memory in the sense that they make use of the Oersted field which is generated by a current passing above (and/or below) an MTJ in dedicated “write lines”. The Oersted field can manipulate the FM layer with the smaller coercive field, the free layer, and does not alter the FM layer with the larger coercive field, the reference layer\textsuperscript{IV}. Depending on the sign of the injected current pulse, and hence of the Oersted field, the relative orientation of the two FM layers can be controlled. The required current amplitude for the Oersted field generation depends on the magnetic properties of the free layer, which can vary strongly between the memory elements due to the fabrication process\textsuperscript{124}. Furthermore, as the device size is reduced the current density injected in the write wires increases in order to achieve the required Oersted field for switching. Hence, field-MRAM showed to be poor in terms of energy consumption and scalability as they are limited to storing elements of about 90 nm diameter\textsuperscript{125}.

\subsection*{2.2.2 Spin-transfer torque switching}

Due to this limitation in scalability, MRAM technology did not take off until the emergence of STT-MRAM in the early 2000s\textsuperscript{15,126}. As shortly discussed in the introductory chapter, a charge current injected perpendicularly into an MTJ gets spin polarized by the fixed reference layer, therefore also called polarizer. This spin current transfers spin angular momentum from the polarizer to the free layer, which is equivalent to a torque acting on the free layer magnetization, the spin-transfer torque (STT)\textsuperscript{8}. From the conservation of spin angular momentum, STT is found to have symmetry $T_{\text{STT}} \propto \mathbf{m} \times (\mathbf{m} \times \mathbf{p})$ with $\mathbf{m}$ being the free layer magnetization and $\mathbf{p}$ the unit vector representing the spin polarization of the injected current, parallel to the direction of the reference layer\textsuperscript{15,126}.

\textsuperscript{IV}In state-of-the-art MTJs the coercive field of the reference layer is increased substantially by pinning it through the exchange interaction to a synthetic antiferromagnet which is fabricated on top of the FM/MO\textsubscript{x}/FM layer.
2.3 Spin-orbit torque switching

The development of STT-MRAM is hampered as several issues remain to be solved \cite{128, 129} of which some we will discuss in the following.

In equilibrium the free and the reference layer of an MTJ are collinear to each other and hence $\mathbf{T}_{STT} = 0$. Hence, thermal agitation, which slightly deviates the magnetization from its equilibrium position, is required to initiate the transfer of angular momentum. The resulting incubation delay can be several $\text{ns}$ and limits the write speed \cite{130, 131}. In principle it can be reduced by e.g. biasing the MTJ with a hard axis field \cite{132}. Additionally, STT-MRAM suffers from the fact that the write and read current paths are identical. This can lead to undesired writing while reading the bit state of small nodes \cite{135, 136}. Furthermore, the RA (product of resistance and MTJ area) has to be small (of the order of $1 \text{Ω} \mu\text{m}^2$) to allow for efficient writing. This requires complex material optimization in order to achieve large enough TMR values for the readout \cite{137}. Finally, the large write currents flowing through the tunnel barrier may result in rapid ageing and reliability issues of the MTJs \cite{136}.

Despite these issues STT-MRAMs are about to be commercialized. Large companies like IBM, Samsung, EverSpin and many more are involved in their development. As an example, EverSpin announced the production of commercially available 256 Mb STT-MRAMs based on 40 nm large perpendicular-MTJs for the end of 2017.

2.3 Spin-orbit torque switching

An alternative to STT is SOT-MRAM which benefits from the in-plane current injection geometry inherent to SOT-induced magnetization switching. In such devices the read and write current paths are decoupled (see Fig. 1.3c) and thereby the electrical stress on the tunnel barrier is strongly reduced. This facilitates the MTJ fabrication substantially and decreases the write error rate due to the separate readout of the MTJ. In the following we give an overview of SOT-induced magnetization dynamics with a focus on current-induced magnetization switching.

As previously mentioned, the demonstration of current-induced magnetization switching in HM/FM/MO$_x$ trilayers with PMA attracted considerable interest. The first experiments were conducted with Pt/Co/AlO$_x$ \cite{29} and Ta-/CoFeB/MgO \cite{98, 138} thin films where switching was observed in the presence of an external in-plane field $\mathbf{B}_x$ for current densities of $10^7 - 10^8 \text{ A cm}^{-2}$. Over the last years successful SOT-induced switching was reported for different HM/FM
2.3. Spin-orbit torque switching

![Graph showing switching probability and critical switching current density as functions of in-plane field and pulse length.]

**Figure 2.5** | **Ultrafast magnetization switching.** a, Switching probability averaged over 100 switching attempts as a function of in-plane field $B_x$ and amplitude of the $\tau_p = 210$ ps long current pulse. With increasing $B_x$ the critical switching current $I_c = I_p (P = 90\%)$ decreases. b, $I_c$ as a function of pulse length $\tau_p$ over 8 order of magnitude. (Figures adapted from Ref. 48)

Consistent with the opposite SOT sign, the switching polarity in Ta/CoFeB/MgO and W/CoFeB/MgO layers is reversed with respect to Pt/Co/AlO$_x$.

Garello et al.\(^{48}\) demonstrated that the magnetization of a 100 nm large Co square can be reversed by injecting sub-ns current pulses. Switching measurements as a function of pulse length $\tau_p$, $I_p$ and $B_x$ showed that, with increasing $B_x$, the injected power required for switching was reduced (see Fig. 2.5a). Furthermore, two distinct switching regimes were evidenced, a thermally activated regime for $\tau_p \gtrsim 1$ ns where the critical switching current density $j_c$ is almost independent of $\tau_p$ and a short-pulse regime ($\tau_p \lesssim 1$ ns) where $j_c$ scales as $\tau_p^{-1}$ (see Fig. 2.5b). Different models have been proposed for the description of the SOT-induced magnetization switching as described in the following sections.

### 2.3.1 Macrospin approximation of switching process

The simplest model is based on the macrospin approximation in which the magnetization is reduced to one magnetic spin described by the unit magnetization

\[^{V}\text{Consistent with the opposite SOT sign, the switching polarity in Ta/CoFeB/MgO and W/CoFeB/MgO layers is reversed with respect to Pt/Co/AlO}_x.\]
2.3. Spin-orbit torque switching

direction vector \( \mathbf{m} = \mathbf{M}/M_s \). Such a model is well suited to describe homogeneous magnetization states with coherent magnetization dynamics (infinite exchange interaction). The dynamic behaviour of such a spin is described by the Landau-Lifshitz-Gilbert (LLG) equation

\[
\frac{d\mathbf{m}}{dt} = -\frac{|\gamma|}{\mu_0(1 + \alpha^2)} \mathbf{m} \times \mathbf{B}_{\text{eff}} - \frac{|\gamma|\alpha}{\mu_0(1 + \alpha^2)} \mathbf{m} \times \mathbf{m} \times \mathbf{B}_{\text{eff}},
\]

(2.14)

with the gyromagnetic ratio of an electron, \( |\gamma| = 1.76 \cdot 10^{11} \text{ rad s}^{-1} \text{T}^{-1} \), the damping constant \( \alpha \) and the saturation magnetization \( M_s \). The effective field \( \mathbf{B}_{\text{eff}} \) is the sum of the magnetic fields acting on the magnetization. Usually, it is the sum of the external field \( \mathbf{B}_{\text{ext}} \) and the effective anisotropy field \( \mathbf{B}_k \) (see Section 2.1). If we consider current induced effects we have to include the SOTs and the Oersted field

\[
\mathbf{B}_{\text{eff}} = \mathbf{B}_{\text{ext}} + \mathbf{B}_k + \mathbf{B}^{FL} + \mathbf{B}^{DL},
\]

(2.15)

Figure 2.6 | Macrospin simulations. Macrospin simulation of SOT-induced magnetization reversal for different combinations of \( B_x, T^{DL} \) and \( T^{FL} \). In the upper row we show the 3 dimensional representation of the magnetization evolution during switching. The lower row shows the corresponding magnetization components, as well as the injected current pulse profile as a function of simulation time. The simulation parameters are: \( a, B_x = 100 \text{ mT}, T^{DL} = 50 \text{ mT}, T^{FL} = 0 \text{ mT}, B_i = 100 \text{ mT} \), \( b, B_x = 0 \text{ mT}, T^{DL} = 50 \text{ mT}, T^{FL} = 0 \text{ mT} \) and \( c, B_x = 100 \text{ mT}, T^{DL} = 50 \text{ mT}, T^{FL} = 32 \text{ mT} \) per \( 10^8 \text{ A cm}^{-2} \).
2.3. Spin-orbit torque switching

Figure 2.7 | DL torque versus external field. 

- **Figure 2.7**

**a**. The simultaneous action of the DL-torque due to the injection of a positive current and \( B_z \) stabilizes the magnetization along \(-z\). **b**. Reversing the current stabilizes the up-state, \(+z\), of the magnetization. (Figure reproduced from Ref. 139)

with \( B^{FL} \) and \( B^{DL} \) being the effective SOT fields defined in Section 2.1.7. The Oersted field is not calculated separately, but, as they have the same symmetry, is regarded as a part of \( B^{FL} \). Depending on the sign of the Rashba-Edelstein effect (and the spin-Hall effect) \( B^{Oe} \) and \( B^{FL} \) add up or are opposed. For Pt/Co/AlO\(_x\) they are opposed.

Alternatively we can realize that \( \mathbf{m} \times \mathbf{B}_{eff} = \sum_i T_i \) and rewrite Eq. (2.14) as

\[
\frac{d\mathbf{m}}{dt} = -\frac{|\gamma|}{\mu_0(1 + \alpha^2)} \sum_i T_i - \frac{|\gamma|\alpha}{\mu_0(1 + \alpha^2)} \mathbf{m} \times \sum_i T_i, \tag{2.16}
\]

with

\[
\sum_i T_i = \mathbf{m} \times \mathbf{B}_{ext} + \mathbf{m} \times [B_k (\mathbf{m} \cdot \mathbf{z}) \mathbf{z}] + T^{FL} + T^{DL}. \tag{2.17}
\]

In either form the LLG-equation can be integrated numerically, e.g. with the Euler-scheme or the Runge-Kutta method, which allows to simulate the SOT-induced magnetization trajectory during current induced switching. In the following we consider the case of a HM/FM sample with an effective out-of-plane anisotropy field \( B_k = 800 \text{ mT} \) and SOTs: \( T^{DL} = 50 \text{ mT} \), \( T^{FL} = 32 \text{ mT} \) per \( 10^8 \text{ A cm}^{-2} \). Starting from the macrospin pointing along \(+z\) we apply an in-plane bias field \( B_x = 100 \text{ mT} \) along the current direction and inject a positive current pulse, in order to switch the magnetization to \(-z\). In Fig. 2.6a-c we show the resulting macrospin trajectories for three particular cases. The lower row shows the magnetization components during the reversal as well as the injected
2.3. Spin-orbit torque switching

Figure 2.8 | Switching phase diagrams. a, Switching phase diagram obtained by macrospin simulation for a $\tau_p = 210\, \text{ps}$ long square current pulse as a function of bias field $B_x$ and pulse amplitude $j_p$ using the parameters $T^{\text{DL}} = 50\, \text{mT}$, $T^{\text{FL}} = 32\, \text{mT}$, and $T^{\text{FL}}_2 = 23\, \text{mT per} \ 10^8\, \text{A cm}^{-2}$ at $T = 0\, \text{K}$. b, Experimentally measured switching phase diagram of a 100 nm large Co square as a function of in-plane field $B_x$ and current amplitude $I_p$ for a constant pulse length of $\tau_p = 210\, \text{ps}$. (Figure reproduced from Ref. 48).

current pulse as a blue solid line, while the upper row is the corresponding three-dimensional representation of the magnetization evolution. Figure 2.6a, b shows the simulation for $T^{\text{FL}} = 0\, \text{mT}$. As expected from experimental reports, the magnetization reverses if an in-plane bias field $B_x = 100\, \text{mT}$ is applied parallel to the current direction (Fig. 2.6a). For $B_x = 0\, \text{mT}$, the switching polarity is not defined because the DL torque does not break the symmetry between the up and down state. Moreover, in the absence of the external field, the energy barrier for reversing the magnetization is larger, which inhibits the switching (Fig. 2.6b). The magnetization reversal symmetry can be described by the simultaneous action of the DL torque and $B_x$ using the simple scheme shown in Fig. 2.7. The external field and $B^{\text{DL}}$, due to a positive current pulse, lead to the magnetization being stabilized along $-z$ where they compensate each other (Fig. 2.7a). Reversing the current pulse destabilizes the down state in favour of the up state (Fig. 2.7b). Hence, $B_x$ is required to break the DL torque symmetry and defines the switching outcome. For negligible FL torque, this simple reasoning can explain the bipolarity of the SOT-induced magnetization switching.

However, as soon as $B^{\text{FL}}$ is taken into account, which acts like a hard-axis field, the macrospin model does not describe the observed switching behaviour correctly any longer. As can be seen in Fig. 2.6c, $B^{\text{FL}}$ induces magnetization
2.3. Spin-orbit torque switching

oscillations about the y-axis, orthogonal to the current direction. Such magnetization dynamics lead to precessional switching, which is characterized by an oscillating switching probability as a function of $B_x$, $j_p$, and $\tau_p$ (Ref. [144][145]). The example in Fig. 2.8 shows a simulated switching phase diagram exhibiting precessional switching as a function of $B_x$ and $j_p$ with both, DL and FL torque taken into account. The experimentally obtained phase diagram of a 100 nm large Pt/Co square illustrated in Fig. 2.8, which does not exhibit any sign of precessional switching, exemplifies the discrepancy of the macrospin model with experiments. In contrast to STT-switching, such oscillatory magnetization switching has not yet been observed in standard SOT-switching experiments [48, 146, 147]. As expected, the macrospin description oversimplifies the reversal process as it neglects micromagnetic effects like the DMI, the exchange interaction $VI$ and the dipolar field. Hence, the macrospin approximation is not suited for describing non-uniform or multi-domain magnetization configurations such as SOT-induced magnetization reversal.

Nevertheless, the macrospin model (ignoring the FL torque), provides a simple analytical expression of the critical switching current [60, 99, 148]. For a sample with PMA, in the limit of $B_x \ll B_k$, the critical switching current density is $^{103}$

$$j_c = \frac{2e M_s t_{FM}}{\hbar} \frac{B_k}{2} \left( B_k - B_x \right), \tag{2.18}$$

with $\xi_{DL}$ defined in Eq. (2.12). Furthermore, the macrospin model correctly describes the action of the SOTs in the small angle approximation of coherent magnetization movement and is used for the analysis of harmonic Hall voltage measurements [97-100] (see Section 3.3).

2.3.2 Switching by domain wall nucleation and propagation

Although, the macrospin approximation can qualitatively reproduce the experimentally obtained switching phase diagrams in certain cases [93], the model has to be improved in order to describe the reversal process correctly in structures larger than the typical domain wall width ($\gtrsim 10$ nm) for which non-coherent magnetization dynamics are expected. In such extended structures the magnetization is best described by more realistic micromagnetic models, in which the system is approximated by an array of macrospins (see Section 2.3.1) interacting $^{VI}$Actually the exchange interaction is included in the macrospin approximation, but it tends to infinity.
with each other by means of exchange interaction, DMI and dipolar field (see Section 2.1).

Micromagnetic simulations, spatially resolved MOKE and all-electrical SOT-induced switching measurements suggest that the reversal occurs through current-induced DW nucleation and propagation. Hence, the magnetization switching is strongly related to current-induced domain wall motion (CIDWM) enabled by the chiral Néel DWs, which were reported in HM/FM/MO\textsubscript{x} trilayers characterized by a strong DMI. In the following we will review the basis of CIDWM and the consequences for SOT-induced magnetization switching. Furthermore, we will cover the different micromagnetic models, which were proposed for the reversal process and relate them to experimental findings.

**Current-induced domain wall propagation**

CIDWM experiments, as well as theoretical predictions and simulations, revealed the importance of the DW structure in enabling the high DW mobilities reported in ultrathin HM/FM bilayers. Starting from the observation of CIDWM, it was suggested that Néel DWs have to be stable in these structures, as Bloch walls cannot be moved efficiently by means of SOTs. (The application of STT to Bloch DWs allowed to move the walls, but in the wrong direction.) From these predictions and experimental results showing that both up-down and down-up DWs move in the same direction for a given current polarity, it was argued that DMI stabilizes chiral Néel DWs. Such DWs are supposed to be very stable and thereby enable CIDWM at high stationary speeds by SOTs. A striking feature is that the DWs move along the current direction (against the electron flow). This is a telltale signature of the SOT-induced DW dynamics and differentiates it from the traditional STT-induced DW displacement, in which the DW is pushed in the direction of the electron flow. Later, experiments, theory, and NV-microscopy directly showed that in Pt/Co/AlO\textsubscript{x} trilayers the DW are left-handed Néel walls illustrated in the lower panel of Fig. 2.2b.

In Fig. 2.9, we illustrate CIDWM in a Pt/Co/AlO\textsubscript{x} trilayer. The DW displacement results from $B_{DL} = B^{DL} \mathbf{y} \times \mathbf{m}$ acting on the central DW moment. The injection of a positive current induces a $B^{DL}$ which pushes the central DW moments of the down-up (up-down) DW towards $-z$ ($+z$). Hence, both DWs propagate at the same speed $v_{DW}$ in the direction of the injected current, against the electron flow. Obviously, the displacement direction will depend on the sign of $B^{DL}$ and the chirality of the DWs. Due to the identical $v_{DW}$
2.3. Spin-orbit torque switching

Figure 2.9 | Current-induced domain wall motion. a, Schematic representation of two left-handed Néel DWs stabilized in a Pt/Co sample due to DMI. The injection of a positive current along $+x$ induces the same DW-velocity $v_{DW}$ for both DWs (yellow horizontal arrows) which leads to a translation of the central domain. CIDWM results mainly from $B^{DL}$ indicated by the two vertical yellow arrows. b, Measurement of $v_{DW}$ as a function of longitudinal bias field $B_x$ and c, transversal bias field $B_y$ for up-down (blue circles) and down-up (red squares) DWs. (Figure adapted from Ref. [50])

of the up-down and down-up DW, the central domain can be pushed back and forth in the wire by reversing the current polarity as the domain size will not be changed. Such a behaviour is required for magnetic memories based on the racetrack technology\cite{59,162,163}.

The application of a moderate in-plane field $B_x$ leads to an asymmetry of $v_{DW}$ for the up-down and down-up DWs (Fig. 2.9b). $B_x > 0$ will increase the domain wall width $\delta_{DW}$ of the down-up DW, while $\delta_{DW}$ of the up-down DW will decrease. This changes their relative response to the SOTs, which leads to distinct $v_{DW}$ for the two DWs\cite{79}. A large $B_x$ can reverse the central DW moments and break the DW chirality induced by the DMI\cite{164}. Due to the reversed chirality, the DWs will propagate in opposite direction and thereby
2.3. Spin-orbit torque switching

expand (compress) the central domain (compare Fig. 2.10b and c). From the critical field $B_{x,c}$, at which the DW starts to move in the opposite direction, the DMI strength can be deduced.

Most often, the effect of the FL torque on CIDWM is ignored in the literature. However, it was shown that the application of a transversal field $B_y$ increases $v_{DW}$ if it points parallel to $B^{FL}$ (see Fig. 2.9c). This is consistent with the results we obtained for SOT-induced switching of Co-dots as a function of bias field $B_y$ shown in Section 4.1.2, which we attribute to a change of $v_{DW}$ with the DW tilt (see below). Finally it was shown that the current-induced domain wall speed can be increased drastically by sophisticated stack engineering. Using two antiferromagnetically coupled wires, current-induced domain wall speeds of up to 750 m s$^{-1}$ per $10^8$ A cm$^{-2}$ were reported. In such structures the effective magnetization is reduced and thus CIDWM does not have to compete with dipolar fields any more.

**Domain wall tilting**

MOKE experiments evidenced that the DW front exhibits a chiral tilt with respect to the propagation direction depending on the current polarity and DW configuration. The DW tilt was shown to be a manifestation of the DMI in such racetracks and was hence proposed to be exploited as a tool to estimate the DMI strength. A detailed discussion of DW tilting is given in Chapter 5.

**2.3.3 Thermally activated switching model**

Two conceptually different micromagnetic models can describe the experimentally observed SOT-induced magnetization switching. The first model that we will describe is based on thermally driven nucleation of reversed domains.

In all the reasoning presented above, temperature effects are neglected, even though the temperature increase due to Joule heating may not be negligible. Hence, Lee et al. developed the thermally activated switching model. We illustrate the effect of the temperature increase due to a current passing in a HM/FM stripe in Fig. 2.10b. The sample is initially saturated along $+z$ (red) and shows thermally activated random nucleation of reversed $-z$ (blue) domains as a result of increasing temperature. Due to the chiral nature of the DWs, $B^{DL}$ will not lead to an expansion of the reversed domains, but rather to a translation of the nucleated domains (see Fig. 2.10b and Section 2.1.6). Therefore, SOT-
induced switching cannot originate from such a random nucleation. However, for in-plane fields $B_x > B_{x,c}$, which break the domain wall chirality dictated by the DMI, an isotropic domain expansion (compression), due to $\mathbf{B}^{DL}$ is expected. Eventually the whole sample will be reversed (see Fig. 2.10c).

In the thermally activated switching model, the required in-plane field, $B_x$, is assumed to be proportional to the effective field stabilizing the chiral DWs. Hence, $B_x$ should be mostly independent of the current amplitude. However, several experimental studies showed that the current amplitude and $B_x$ are strongly coupled\cite{64, 68, 69}. If either is decreased the other has to increase (see Fig. 2.5). This observation contrasts with the thermally activated switching model described above. Furthermore, micromagnetic simulations showed that the switching process is robust with respect to temperature increase in Pt/Co/AlO$_x$ thin films with strong PMA.

On the other hand, time-resolved MOKE measurements of SOT-switching in

Figure 2.10 | Thermally activated switching. a, Micromagnetic simulation of thermally activated random nucleation of down (blue) domains in a stripe which was initially saturated along $+z$ (red). (Figure adapted from Ref. [64]) b, Injection of a positive current leads to a translation of the up-down and down-up DW at the same speed $v_{DW}$ and in the same direction as long as $B_x < B_{x,c}$ and hence the domain size is conserved during CIDWM. c, Application of an in-plane bias field $B_x > B_{x,c}$ breaks the DMI induced chirality of the up-down DW, and hence the down domain expands as a result of the reversed $v_{DW}$ for the up-down DW.
2.3. Spin-orbit torque switching

$1\,\mu\text{m}$ large Pt/Co squares showed a reproducible homogeneous switching, which is accompanied by a slow recovery of the magnetization after the current pulse. This after pulse dynamics are probably caused by a temperature increase during current pulse injection\cite{62}. Similar results were obtained by Yoon et al. on $3\,\mu\text{m}$ large Ta/CoFeB squares. However, they relate the slow after-pulse relaxation observed by MOKE to FL torque promoted DW reflections at the sample edge rather than to a temperature increase\cite{168}. In order to rule out heating effects the sample size can be reduced, as for a given switching current density $j_c$ and resistivity $\rho$, the injected power, and hence Joule heating is proportional to the width and the length of the current line. However, in order to resolve the switching process in such small samples high spatial resolution is required.

2.3.4 Edge-driven nucleation

Pizzini et al. demonstrated in $70\,\mu\text{m}$ wide Pt/Co/AlO$_x$ structures that DMI plays a crucial role in determining the DW nucleation point\cite{150}. In their experiment they applied various in-plane fields $B_x$ and measured the response of the FM to magnetic field pulses $\pm B_z$ of constant amplitude. Figure 2.11a-d summarizes the results they obtained by MOKE microscopy. At $B_x = 0\,\text{mT}$ the application of $-B_z$ was not sufficient to induce any DW nucleation in the initially along $+z$ saturated sample (Fig. 2.11a). At $B_x = 260\,\text{mT}$ they observed DW nucleation on the left sample edge due to $-B_z$ pulses (Fig. 2.11b). Reversing $B_x$ changed the nucleation side (Fig. 2.11c) and under $B_x = 215\,\text{mT}$, an initially along $-z$ saturated sample, showed DW nucleation on the right side due to $+B_z$ pulses. This peculiar DW nucleation symmetry could be attributed to the DMI, as we illustrate in Fig. 2.11d. In finite structures DMI tilts the equilibrium edge magnetization symmetrically inwards (outwards) for $+M_z (-M_z)$. The in-plane field $B_x$ breaks the symmetry of the spin canting (Fig. 2.11f) and favours DW nucleation due to $B_z$ on the side with larger $M_x$-component (Fig. 2.11g).

Figure 2.12a shows micromagnetic simulations of field-induced magnetization reversal in a $90\,\text{nm}$ Co-square by Martinez et al.\cite{66}. In agreement with the results obtained by Pizzini et al. in large structures, Martinez et al. observed magnetization switching through a DW nucleation and propagation process.

Different groups realized that a similar reasoning can be applied to SOT-induced magnetization switching in Pt/Co/AlO$_x$ samples with the Co layer patterned into a finite structure, e.g. a circular Co dot on top of a Pt current line\cite{65,66}. Replacing the $\pm B_z$ field pulse by $\mathbf{B}^{DL}$ due to a current in the Pt line, turned out to give similar results. Figure 2.12b shows the micromagnetic
2.3. Spin-orbit torque switching

![Image](image_url)

**Figure 2.11 | DMI driven edge nucleation.** a, MOKE images showing the initial state \((+M_z)\) of the 70 µm wide Pt/Co structures. b, Under a constant \(B_x = 260\) mT a field pulse \(-B_z\) is applied. DW nucleation is observed on the left side of the structure indicated by the white arrow. c, With \(B_x = -260\) mT the DW nucleates on the right side of the structure. d, Starting from a negative initial state and \(B_x = 215\) mT a DW nucleates on the right side of the structure if a \(+B_z\) field pulse is applied. (Figures adapted from Ref. 150). e, In equilibrium the edge magnetization is tilted inwards (outwards) for \(+M_z\) \((-M_z)\) due to DMI. f, \(B_z\) breaks the symmetry of the edge canting which defines the favoured nucleation side. g, \(B_z\) reverses the magnetization at the edge with the largest \(M_x\) component and pushes the nucleated DW through the sample.

Simulation by Mikuszeit et al.\(^{65}\) of a 100 nm wide Co dot which, indeed, shows that the magnetization reverses through edge-driven DW nucleation followed by a current-induced DW propagation process. In Fig. 2.12, we show the corresponding evolution of the normalized \(M_z\) component during reversal for different current densities \(j_p\), given in units of A m\(^{-2}\). As soon as the critical switching current density \(j_c\) is overcome (blue line) the switching time decreases with increasing \(j_p\). At very high \(j_p\) (green line) turbulent magnetization dynamics were observed. Such micromagnetic simulations conducted at finite temperatures showed that the switching process is very robust with respect to temperature and can nicely reproduce the results obtained in all-electrical transport measurements for Pt/Co/AlO\(_x\) based structures\(^{48}\). However, for devices based on Ta/CoFeB/MgO, this model did not succeed yet in reproducing the experimen-
2.3. Spin-orbit torque switching

Figure 2.12 | Micromagnetic simulations of edge driven nucleation. a, Micromagnetic simulation of field-induced magnetization reversal in a 90 nm large Co square by Martinez et al.\textsuperscript{[66]} b, Micromagnetic simulation of current-induced magnetization switching of a circular 100 nm wide Co dot. The magnetization reverses through a current-induced DW nucleation and propagation process in agreement with the reasoning presented by Pizzini et al.\textsuperscript{[150]}. b, Evolution of the normalized $M_z$ component during the current pulse for different current densities given in A m$^{-2}$.  

...tally observed switching behaviour\textsuperscript{[30]} using reasonable parameters. Compared to Pt(3 nm)/Co(0.6 nm)/AlO$_x$, which has a DMI of $\approx 2$ mJ m$^{-2}$, Ta/CoFeB/MgO has typically a negligible DMI\textsuperscript{[169]}, which leads to stable Bloch walls in these bilayers\textsuperscript{[161]}.

The aim of the remaining part of this thesis is to shed light on the microscopic mechanisms and time-resolved dynamics leading to SOT-induced magnetization switching in 500 nm wide Co(1 nm) dots on a Pt(5 nm) current-line. The magnetization reversal process is studied by time- and spatially-resolved scanning transmission x-ray microscopy and related to all-electrical switching experiments and micromagnetic simulations. As to date, except the time-resolved MOKE-microscopy measurements introduced previously, which have a spatial resolution of several hundred nm, neither time- nor spatially-resolved measurements of SOT-induced magnetization switching were reported and several questions of fundamental as well as of technical interest remain open.
3 Methods

3.1 Sample fabrication

In the following sections we discuss the fabrication processes for the preparation of the various samples used in this thesis. All the explanations are based on Pt/Co/AlO$_x$ samples with PMA. First, we focus on the thin film deposition, as it is the basis for the different fabrications processes described in the subsequent sections.

3.1.1 Thin-film deposition

We deposit the Pt/Co/AlO$_x$ thin films by dc magnetron sputtering at a base pressure of $4 \cdot 10^{-8}$ mTorr, with the substrate rotating at 30 rpm. Table 3.1 summarizes the typical process parameters used for the growth of Pt(5 nm)/Co(1 nm)/AlO$_x$ films. Prior to the thin film deposition the Si/SiO$_2$ (Si/Si$_3$N$_4$) substrates are cleaned by means of Ar$^+$ milling. Thereafter, the Pt/Co/Al trilayer is deposited, followed by an in situ Al oxidation step in an oxygen plasma. As the effective out-of-plane anisotropy field ($B_k$) depends strongly on the oxygen content at the Co/AlO$_x$ interface, we can tune $B_k$ precisely by adjusting the oxidation time. The parameters summarized in Table 3.1 result in $B_k \approx 680$ mT (compare Fig. 3.8) which gives a good compromise between PMA and SOT-induced magnetization switching at reasonable current densities (see Section 4.1).

<table>
<thead>
<tr>
<th>Process</th>
<th>Power (W)</th>
<th>Current (mA)</th>
<th>Pressure (mTorr)</th>
<th>Process time (s)</th>
<th>Thickness (nm)</th>
</tr>
</thead>
<tbody>
<tr>
<td>Ar$^+$ cleaning</td>
<td>50</td>
<td>250</td>
<td>10</td>
<td>60</td>
<td>-</td>
</tr>
<tr>
<td>Pt deposition</td>
<td>7</td>
<td>20</td>
<td>0.5</td>
<td>52</td>
<td>5</td>
</tr>
<tr>
<td>Co deposition</td>
<td>9</td>
<td>25</td>
<td>0.6</td>
<td>25</td>
<td>1</td>
</tr>
<tr>
<td>Al deposition</td>
<td>15</td>
<td>40</td>
<td>1.2</td>
<td>65</td>
<td>1.6</td>
</tr>
<tr>
<td>Al oxidation</td>
<td>25</td>
<td>190</td>
<td>13.2</td>
<td>35</td>
<td>-</td>
</tr>
</tbody>
</table>

Table 3.1 | Process parameters for the deposition of Pt(5 nm)/Co(1 nm)/AlO$_x$ thin films by dc magnetron sputtering in chronological order from top to bottom. The Al oxidation time is optimized for sample with $B_k \approx 680$ mT.
3.1. Sample fabrication

3.1.2 Fabrication of Hall bars

The section below sketches the fabrication processes of Hall bars using optical lithography. We follow two different approaches. The first is based on Ar$^+$ milling, while the second uses a lift-off process.

**Ion milling process**

After the deposition of the thin films described in the previous paragraph (step 1, 2 in Fig. 3.1a), we can fabricate micrometer sized structures (i.e. Hall bars) by means of Ar$^+$ milling. Such a process is sketched in Fig. 3.1a. We remove...
potential adsorbates on the film surface (e.g. water), by cleaning the sample thoroughly in a 1 mBar O\textsubscript{2}-plasma for 90 s at a power of 150 W followed by subsequently cleaning in acetone and isopropyl alcohol (IPA) at 50°C for 300 s. This is the standard cleaning procedure in the ensuing fabrication steps. Thereafter, we spin coat the sample with ma-N 1420 (negative resist, 2 \( \mu \)m thick) using a one step ramp: 0 rpm to 4000 rpm in 4 s, 4000 rpm for 45 s (step 3). After spinning, the resist is pre-baked for 125 s at 120°C. Step 4 shows how the sample is illuminated in the MJB3 maskaligner with UV radiation of 10 mW cm\(^{-2}\) intensity and 310 W for 35 s through an optical mask. The polymer of the negative ma-N resist is crosslinked by the UV radiation and therefore, the parts which have been exposed are no longer soluble in the developer (ma-D 553s). Step 5 shows the sample after rinsing it for 42 s in ma-D 553s and a couple of minutes in deionized water. The resist, which has not been illuminated, is completely removed by the developer. Care has to be taken as most of the developers are bases, which attack the AlO\textsubscript{x} capping layer and may cause damage to the sample. Finally, we fabricate the Hall bar by means of an Ar\textsuperscript{+} milling (step 6) process using the resist as a milling mask. The Ar\textsuperscript{+} milling is done with 390 W (50 mA) at a pressure of 0.18 mTorr in a dedicated ion miller. Thereafter, the remaining resist can be removed with the standard cleaning procedure. Figure 3.2a shows the negative optical mask used in step 4.

### Lift-off process

In contrast to the ion milling process where the first step is the thin film deposition, in a lift-off process the structure is defined before the deposition (see Fig. 3.1b). First, the optical resist ma-N 1420 is spun on the cleaned bare substrate with the same settings as in the previous paragraph. In a lift-off process the resist is not used as a milling mask, therefore, we harden the resist at a reduced temperature of 110°C for 125 s to facilitate the resist removal and lift-off. After the pre-bake, the sample is illuminated through an inverted optical mask (see Fig. 3.2a and 3.2b) by UV radiation (10 mW cm\(^{-2}\), 310 W, 35 s) and developed in ma-D 553s for 50 s. The reduction of the pre-bake temperature and the increase of the development time leads to an undercut in the resist (see step 4), which avoids vertical edges during the deposition and thereby enables lift-off. Once the structure is defined by the inverted resist mask, the thin film is deposited onto the substrate (see Section 3.1.1). Due to the resist, the Ar\textsuperscript{+}

\(^{1}\)The O\textsubscript{2}-plasma does not only clean the surface, but it also increases the wetting of the resist during spin coating.
3.1. Sample fabrication

Figure 3.2 | **Masks for optical lithography.** a, Negative optical mask used for Ar$^+$ milling based processes and b, positive optical mask used for lift-off based processes. The UV radiation can pass through the white area and does not pass the Cr (yellow). c, Scanning electron micrograph of a 10$\mu$m wide Hall bar fabricated by lift-off.

Cleaning step during the deposition process is omitted as the organic resist may contaminate the sputtering chamber otherwise. After deposition, the sample is left in acetone for 12 hours at room temperature. This removes the exposed resist and the deposited layers on top of it (step 7). In case the lift-off does not work, the resist removal can be facilitated by blasting it off with a syringe filled with acetone. Finally, the sample is cleaned with the standard procedure in order to remove any resist residues on the substrate surface. Figure 3.2 shows a scanning electron micrograph of the resulting Hall bar.

3.1.3 Fabrication of magnetic dots

**Electron beam lithography**

Additionally to the optical lithography described in the two previous paragraphs, we use electron beam lithography (EBL) for the fabrication of sub $\mu$m-sized structures. Instead of the negative optical resist ma-N 1420 we use the positive EBL resist PMMA 950k 1:1 EL. We spin the resist on the cleaned sample with a two step ramp: 0 rpm to 1000 rpm in 1 s, 1000 rpm for 4 s, 1000 rpm to 3000 rpm in 4 s, 3000 rpm for 45 s. The samples have to be pre-baked for 600 s at 180°C in order for the solvent to evaporate. This results in crosslinked PMMA of around 200 nm thickness. According to the EBL design, the PMMA layer is locally destroyed by the focused electron beam of the EBL. During development in MIBK 1:3 for 60 s the damaged PMMA is removed. In order to stop the
3.1. Sample fabrication

In the following, the fabrication process for magnetic dots is described (compare Fig. 3.3). The same processes can be used for different material compositions.
and substrates, by slightly adapting some of the fabrication steps.

**Step 1 - 2** Deposition of the Pt/Co/AlO\(_x\) by dc magnetron sputtering (see Section 3.1.1).

**Step 3 - 7** Fabrication of the Au contacts by optical lithography using a lift-off process (see Section 3.1.2).

**Step 8 - 12** Fabrication of a Ti dot defined by EBL using a lift-off process. The Ti dot is used as a hard mask for the ion milling, similar to the resist mask in step 5 of the ion-milling process.

**Step 13** Ion milling of the Co/AlO\(_x\) with optimized milling time in order to stop at the Pt/Co interface (see Section 3.1.2). The thickness of the Ti dot deposited during step 11 should be sufficient to protect the Co/AlO\(_x\) beneath, however, it should also not be too thick, as otherwise the current we will shunt through the Ti. We used 7 nm Ti.

**Step 14 - 18** Fabrication of a Ti current line (or Hall cross) by EBL using a lift-off process. The Ti line is used as a hard mask for the milling.

**Step 19** Definition of the current line by Ion milling of the Pt. Similar to step 13 the Ti thickness deposited during step 17 should be sufficient to work as a hard mask, but as thin as possible for minimizing current shunting. Typical Ti thickness for a Pt(5 nm) line is 5 nm.

Figure 4.1h shows an atomic force micrograph of a typical 500 nm large Ti-capped Co dot patterned on top of a Hall cross.

### 3.2 Electrical measurements

In this section we provide the details about the electrical measurements conducted within the scope of this thesis. This includes the circuitry, the signal generation and acquisition, as well as typical measurement schemes. The section is divided in two parts; the first part focuses on ac harmonic Hall voltage measurements, while the second part focuses on pulsed magnetization switching experiments.

If not stated differently the following instruments have been used:

- Current source: Keithly 6221 dc- and ac-current source
3.2. Electrical measurements

- Voltage source: NI PXI-4461 in NI PXI-1036DC chassis
- Pulse generator: RTV40 sub-ns pulser
- Electromagnet: GMW 5403 delivering up to 1.8 T
- Sampling oscilloscope: Infiniium DSO81304A (13 GHz, 40 Gsa/s)
- Signal acquisition: NI PXI-4462 in NI PXI-1036DC chassis
- Instrument control with LabVIEW 2016

3.2.1 Harmonic ac Hall voltage measurements

Figure 3.4 shows the circuitry used for the standard harmonic ac Hall voltage measurements. This measurement scheme is used for quasi static magnetization characterization and the measurement of current induced effects.

In such a measurement scheme a low frequency \( (\omega = 2\pi \cdot 10 \text{ Hz}) \) ac-current, typically resulting in a current density in the sample of \( j = 1 \cdot 10^7 \text{ A cm}^{-2} \), is injected laterally along the \( x \)-direction into a Hall bar. The resulting transversal and longitudinal voltages \( V_H \) and \( V_{xx} \) are acquired by the PXI measurement system and fast Fourier transformed by a customized LabVIEW program in order to access the first and second harmonic components of the signals. The sample is mounted in the electromagnet on a motorized stage, which allows to align it at any angle with respect to the magnetic field of up to 1.8 T. In a typical field scan the azimuthal \( (\varphi_B) \) and polar \( (\theta_B) \) angle of the magnetic field \( B_{ext} \) are

\[ \begin{align*}
V_H & \quad \text{Transversal voltage} \\
V_{xx} & \quad \text{Longitudinal voltage} \\
\end{align*} \]

**Figure 3.4 | Electrical measurement setup.** a, Circuitry for the adiabatic harmonic Hall voltage measurements: An ac-current of frequency \( \omega = 2\pi \cdot 10 \text{ Hz} \) is injected laterally into the sample along the \( x \)-direction and the resulting transverse Hall voltage \( (V_H) \) as well as the longitudinal voltage \( (V_{xx}) \) are measured. b, Illustration of the coordinate system: \( x \) is defined to be parallel to the direction of a positive (technical) current and \( z \) points along the sample normal. The magnetic field \( B_{ext} \) is defined by the azimuthal \( (\varphi_B) \) and polar \( (\theta_B) \) angle.
3.2. Electrical measurements

fixed (see Fig. 3.4b), while the magnetic field is swept from positive to negative and back to its initial state. At each field point \( V_H \) and \( V_{xx} \) are acquired.

### 3.2.2 Pulsed magnetization switching measurements

**Hall voltage based**

Figure 3.5a shows the circuitry for the all-electrical pulsed magnetization switching measurements based on the Hall voltage. Fast voltage pulses with amplitudes \( U_p \) of up to 35 V and pulse lengths \( \tau_p \) ranging between 250 ps and 20 ns are generated by the pulser and injected laterally into the sample. We read the magnetization state after each pulse with a small ac-voltage (typically 20-200 mV amplitude, \( \omega = 2\pi \cdot 635 \text{ Hz} \)) applied by the PXI-4461 card and measure the resulting Hall voltage with the PXI-4462 module. The increased frequency, compared to the harmonic ac Hall voltage measurements discussed in Section 3.2.1, enables the data acquisition with 10 ms integration time by using a lock-in amplifier based analysis. Simple multiplication of \( V_{100} \) (voltage drop on a 100 \( \Omega \) resistor connected in series with the sample) with the normalized reference signal \( (V_{ref}) \) allows to determine the injected current. Multiplication of \( V_{ref} \) with

![Pulsed switching setup](image)

**Figure 3.5 | Pulsed switching setup.** a, Circuitry for SOT-induced magnetization switching measurements: Fast voltage pulses are injected laterally into the sample and monitored on a fast sampling oscilloscope using a 20 dB pick-off tee. The magnetization state is measured after each pulse by means of transversal Hall voltage \( (V_H) \) due to the application of a low amplitude ac voltage with frequency \( \omega = 2\pi \cdot 635 \text{ Hz} \). b, Circuitry used for SOT-induced magnetization switching of MTJ’s. The MTJ state is measured by means of TMR after each pulse. In order to prevent the fast voltage pulses to go through the junction, the read and write paths are separated by a bias tee.
3.2. Electrical measurements

$V_H$ permits to distinguish between the two remanent magnetization states ($+M_z$ and $-M_z$) and the quantification of the Hall effect by normalizing it by the injected current amplitude (see Section 3.3). We separate the fast pulses from the read path by a bias tee. Furthermore, we impedance match the sample by connecting $50 - 62 \, \Omega$ SMD-resistors in parallel and close to the high resistance sample and reduce pulse reflections by mounting the sample on a $50 \, \Omega$ coplanar wave guide. Finally, we try to avoid spreading of the voltage pulses into the Hall contacts by adding $100 \, k\Omega$ SMD-resistors in series with the Hall contacts.

We make use of a $20 \, \text{dB}$ pick-off tee to monitor the injected current pulses with a fast sampling oscilloscope. The signal-to-noise ratio may be improved in the future by replacing the $100 \, k\Omega$ SMD-resistors with SMD-inductors having a cut-off frequency adapted to the pulse frequency. The sample is mounted in the electromagnet, which allows to apply a magnetic field in well defined directions.

A typical all-electrical SOT-induced switching measurement follows the procedure described by Garello et al. During such a measurement, the external field is kept constant with $\varphi_B = 0^\circ$ and $\theta_B = 89^\circ$, almost parallel to the current line, and the effect of single voltage pulses with well defined length ($\tau_p$) and amplitude ($U_p$) on the magnetization is measured. In order to do so, we have to ensure the initial magnetization state before injection of the write pulse. This is achieved by defining a reset pulse with amplitude and duration optimized to reset the magnetization (typically $\tau_p = 20 \, \text{ns}$) to the desired initial state ($+M_z$ or $-M_z$). We make sure that the magnetization was correctly reset by measuring $R_H$ after each reset pulse and, if necessary, repeat the reset several times. Once the magnetization state is reset correctly, we apply a single set pulse of opposite polarity and measure the difference of $R_H$ before and after the injection. If $\Delta R_H$ corresponds to $R_{AHE}$ the magnetization switching was successful for the injected set pulse. We average over up to 1000 switching attempts per $\tau_p$ and $U_p$ in order to determine the switching probability for a given set of parameters. If the sample has stable intermediate magnetization states (i.e. multi-domain states) we can average $\Delta R_H$ and obtain a signal which is proportional to the fraction of the dot that reverses.

**Tunneling magneto resistance based**

A similar setup can be used for the all-electrical switching experiments in MTJs. In such measurements the magnetization state is detected by exploiting the TMR (see introductory chapter). As the tunnel barriers are fragile we have to ensure that the in-plane (write) voltage pulses are restricted to the current line and do
3.3. Harmonic Hall voltage measurements

not escape through the junction, potentially damaging it. For this purpose we
connect the dc-port of a bias tee to the MTJ’s top contact (as shown in Fig. 3.5b)
through which we read the MTJ state with the lock-in amplifier based analy-
sis method described in the previous paragraph. Typically the measurement
procedure follows the set-reset pulse approach describe above.

Additionally, the separation of the fast write and the slow read path through
the bias tee enables the application of constant dc STT-biases during the injec-
tion of the fast current pulses.

3.3 Harmonic Hall voltage measurements

In HM/FM/MOₓ samples, the SOTs can be quan-
tified as a function of magnetization direction with
high accuracy by conducting adiabatic Hall voltage measurements.

Injecting an in-plane ac current

\[ I(t) = I_0 \cos(\omega t) \]

with frequency \( \omega = 2\pi \times 10 \text{ Hz} \) into the sample in the \( x \)-direction leads to the generation
of a Hall voltage \( V_H = V_{xy} \) perpendicular to the cur-
rent direction, along the \( y \)-direction (see Fig. 3.4). \( V_H \)
depends on the magnetization direction through the
planar and anomalous Hall effect (PHE, AHE) and is given by

\[
V_H = I \left[ R_{AHE} \cos(\theta) + R_{PHE} \sin^2(\theta) \sin(2\varphi) \right],
\]

with \( R_{AHE} \) and \( R_{PHE} \) being the anomalous and the planar Hall coefficients, respectively. \( \theta \) and \( \varphi \) are the polar and azimuthal magnetization angles, re-
spectively, defined in Fig. 3.6. The injected ac current induces oscillating SOTs
which in turn lead to the magnetization oscillating around the equilibrium po-

tion. Hence, the transversal voltage changes as a function of time as given
by

\[
V_H(t) = I(t) \left\{ R_H \left[ B_0 + B^I(t) \right] \right\}.
\]

The magnetic field \( B_0 \) is the sum of the time independent external field \( B_{\text{ext}} \) and
the effective anisotropy field (magnetocrystalline anisotropy and demagnetizing
field) \( B_k \). \( B^I(t) \) parametrizes the current induced effective fields (SOTs and
Oersted field), which are expected to be directly proportional to \( I(t) \). Therefore,

\[^{1} \text{We neglect the ordinary Hall effect which would appear as an additional term } (R_{OHE} B_{\text{ext}} \cos(\theta_B)) \text{ in Eq. (3.1). This is reasonable, as } R_{OHE} \text{ is negligibly small compared to } R_{AHE} \text{ and } R_{PHE} \text{ in ferromagnetic samples.}\]
3.3. Harmonic Hall voltage measurements

Eq. (3.2) can be approximated as

\[ V_H(t) \approx I(t) \left[ R_H(B_0) + \frac{dR_H}{dB} \cdot B \cos(\omega t) \right]. \]  

(3.3)

In terms of total resistance, \( R_H = V_H/I \), the ac current generates a signal with a first and second harmonic component: \( R_H = R_H^{\omega} + R_H^{2\omega} \) with \( R_H^{\omega} = R_H(B_0) \) and \( R_H^{2\omega} = \frac{1}{2} \frac{dR_H}{dB} \cdot B \).

The detailed calculations can be found in Appendix A.3.

3.3.1 First harmonic effects

The first harmonic component \( R_H^{\omega} \) (Eq. (3.1) normalized by the current amplitude \( I_0 \)) is linked to the equilibrium magnetization position given by \( B_{ext} \) and \( B_k \). Figures 3.7 a-c show the first harmonic signals, obtained by field sweep measurements with \( \varphi_B = 0^\circ, 45^\circ, 90^\circ \) and fixed \( \theta_B = 86^\circ \) according to the measurement procedure introduced in Section 3.2.1. Obviously, the \( R_H^{\omega}(\varphi_B = 0^\circ, 90^\circ) \) curves are both purely antisymmetric with respect to \( B_{ext} \). This is the case, as in a sample with uniaxial out-of-plane anisotropy \( \varphi_B = \varphi \) holds, which leads to a vanishing \( R_{PHE} \) term in Eq. (3.1). Accordingly, \( R_H^{\omega}(\varphi_B = 45^\circ) \) has an

Figure 3.7 | Harmonic Hall voltage measurements of SOTs. a-c, First harmonic Hall resistance signals obtained for \( \varphi_B = 0^\circ, 45^\circ \) and \( 90^\circ \) at constant \( \theta_B = 86^\circ \). d-f, Corresponding second harmonic signals.
3.3. Harmonic Hall voltage measurements

Figure 3.8 | Analysis of first harmonic signal. a, First harmonic signal measured at \( \varphi_B = 45^\circ \) and \( \theta_B = 86^\circ \). b, Decomposition of \( R_H^o(\varphi_B = 45^\circ) \) into its symmetric \( (R_H^o(\text{sym.}) = R_{PHE} \sin^2(\theta) \sin(2\varphi)) \) and antisymmetric \( (R_H^o(\text{asym.}) = R_{AHE} \cos(\theta)) \) component. c, Evaluation of \( \theta \) as a function of \( B_{\text{ext}} \) using the antisymmetric component. d, Determining \( R_{PHE} \) from the symmetric component. e, Evaluation of the effective anisotropy field as a function of \( \sin^2(\theta) \). The two dashed lines indicate the error in the analysis of \( B_k \) due to a misalignment of \( \theta_B = 86 \pm 1^\circ \).

additional symmetric component. Therefore, the separation of \( R_H^o(\varphi_B = 45^\circ) \) into the symmetric and antisymmetric parts allows to quantify \( R_{AHE} \) and \( R_{PHE} \) (see Figs. 3.8a and b).

The antisymmetric part of the signal \( (R_H^o(\text{asym.}) = R_{AHE} \cos(\theta)) \) allows to evaluate \( \theta \) for each \( B_{\text{ext}} \) (see Fig. 3.8c) through

\[
\theta(B_{\text{ext}}) = \cos^{-1} \left( \frac{R_B^o(B_{\text{ext}})}{R_{AHE}} \right).
\]

The anomalous Hall coefficient can easily be deduced from a field scan, as \( R_{AHE} = R_H^o(B_{\text{ext}} = 0) \) in the case of a sample with PMA. Contrary, \( R_{PHE} \) is obtained by analysing the symmetric component of \( R_H^o(\varphi_B = 45^\circ) \). It is simply the slope of a linear fit to a plot of \( \sin^2(\theta) \) versus \( R_H^o(\text{sym.}) \) (compare the last term in Eq. (3.1) and Fig. 3.8d).

As the first harmonic signal is linked to the equilibrium position of the magnetization, it is possible to determine \( B_k \) from such a measurement. The equilibrium position of the magnetization is defined by the simultaneous action of \( B_{\text{ext}} \) and \( B_k \). Hence, in equilibrium the condition \( T_Z + T_k = 0 \) has to hold. \( T_Z \) is the Zeeman torque originating from the external field and \( T_k \) is the magnetic torque from the effective anisotropy field. After a short calculation using \( \varphi = \varphi_B \) for
out-of-plane magnetized samples, we find

\[ B_k = B_{\text{ext}} \left[ \frac{\sin(\theta_B)}{\sin(\theta)} - \frac{\cos(\theta_B)}{\cos(\theta)} \right] . \]  

(3.5)

\( B_k \) is calculated for every \( B_{\text{ext}} \) value by replacing \( \theta \) with \( \theta(B_{\text{ext}}) \) from Eq. (3.4). This allows us to identify possible angle dependences of \( B_k \) by plotting it as a function of \( \sin^2(\theta) \) (see Fig. 3.8). The detailed calculation can be found in Appendix A.1. As can be seen in Fig. 3.8, the analysis fails for large values of \( \sin^2(\theta) \). This has its origin in small errors during the \( \theta_B \) calibration, which are amplified by the analysis for large values of \( \sin^2(\theta) \). As example the dashed lines in Fig. 3.8 indicate the resulting \( B_k \) by assuming a misalignment of \( \theta_B = 86 \pm 1^\circ \).

### 3.3.2 Second harmonic effects

The second harmonic component \( R^{2\omega}_H = \frac{1}{2} \frac{dR_H}{dB_I^l} \cdot B^l \) is connected to current induced effective fields originating from the SOTs and the Oersted field III. If we calculate the above derivative by replacing \( R_H \) with the corresponding first harmonic signal \( R^\omega_H \) and keep in mind that \( \theta(B^l) \) and \( \varphi(B^l) \) we find

\[
R^{2\omega}_H = \frac{1}{2} \frac{dR_H}{dB_I^l} \cdot B^l = \frac{1}{2} \left[ R_{\text{AHE}} - 2R_{\text{PHE}} \cos(\varphi) \sin(2\varphi) \right] \frac{d\cos(\theta)}{dB_I^l} \cdot B^l \\
+ \frac{1}{2} R_{\text{PHE}} \sin^2(\theta) \frac{d\sin(2\varphi)}{dB_I^l} \cdot B^l.
\]  

(3.6)

For a field scan (constant \( \theta_B \) and \( \varphi_B \)) the derivatives appearing in Eq. (3.6) can be approximated in spherical coordinates with the small angle approximation to

\[
\frac{d\cos(\theta)}{dB_I^l} \cdot B^l = \frac{d\cos(\theta)}{dB_{\text{ext}}} \frac{1}{\sin(\theta_B - \theta)} B^l_\theta,
\]  

(3.7)

and

\[
\frac{d\sin(2\varphi)}{dB_I^l} \cdot B^l = \frac{2 \cos(2\varphi)}{B_{\text{ext}} \sin(\theta_B)} B^l_\varphi.
\]  

(3.8)

We do consider only the effect of the polar \( (B^l_\theta) \) and azimuthal \( (B^l_\varphi) \) components of the current induced effective fields as the radial component does not affect the motion of the magnetization and is therefore ignored in the discussion. We use the definition of the FL SOT \( (T^{FL} = T^{FL} \mathbf{m} \times \mathbf{y}) \) and the DL SOT \( (T^{DL} = \)\textit{We do not treat the Oersted field separately in the following, as it can not be distinguished experimentally form the FL torque due to the identical symmetry of both effects. The extracted FL torque has to be corrected by the Oersted field contribution using model calculations.}
3.3. Harmonic Hall voltage measurements

\[ T^{DL} \mathbf{m} \times (y \times \mathbf{m}) \] given in Section 2.1.7 and consider the corresponding effective fields \( \mathbf{B}^{FL} = B^{FL}y \) and \( \mathbf{B}^{DL} = B^{DL}y \times \mathbf{m}, \) respectively. In spherical coordinates these effective fields can be decomposed into their polar \( (B_\theta = \mathbf{B} \cdot \mathbf{e}_\theta) \) and azimuthal \( (B_\varphi = \mathbf{B} \cdot \mathbf{e}_\varphi) \) components by carrying out the scalar product with the corresponding unit vectors in spherical coordinates \( \mathbf{e}_\theta(\theta) \) and \( \mathbf{e}_\varphi(\varphi) \) (compare Fig. 3.6).

\[ B_I^{\theta}=B_{DL}^{\theta}+B_{FL}^{\theta} \quad \text{and} \quad B_I^{\varphi}=B_{DL}^{\varphi}+B_{FL}^{\varphi}, \] respectively. Following this reasoning we rewrite Eq. (3.6) as

\[
R_2^\omega \propto \frac{1}{2} \left[ R_{AHE} - 2 R_{PHE} \cos(\theta) \sin(2 \varphi) \right] \frac{d \cos(\theta)}{dB_{ext}} \frac{1}{\sin(\theta_B - \theta)}
+ \frac{1}{2} R_{PHE} \sin^2(\theta) \frac{2 \cos(2 \varphi)}{B_{ext} \sin(\theta_B)} \left[ B_{FL} \cos(\varphi) - B_{DL} \sin(\varphi) \cos(\theta) \right] + R_{VT}^{2\omega}.
\] (3.9)

The last term in Eq. (3.9) is due to the thermoelectric voltage, which is mainly driven by the anomalous Nernst effect (ANE) induced by Joule heating. This thermoelectric voltage is proportional to \( \nabla T \times \mathbf{m} \) where \( \nabla T \) is the temperature gradient induced by the injected current. The heat dissipation is mainly through the substrate as, compared to air, it is a good heat conductor. This leads to the temperature gradient mainly perpendicular to the sample plane along the \( z \)-direction which induces a thermoelectric voltage in \( y \)-direction

\[
V_{VT} \propto \nabla T \cos(\varphi) \sin(\theta).
\] (3.10)

As the temperature gradient is proportional to the power injected into the sample and hence \( \nabla T \propto I(t)^2 \propto I_0^2 [1 + \cos(2\omega t)] \), the second harmonic signal

\[
R_{VT}^{2\omega} = I_0 \nabla T \sin(\theta) \cos(\varphi),
\] (3.11)

is generated. During the SOT analysis \( R_{HT}^{2\omega} \) has to be corrected by the \( R_{VT}^{2\omega} \) term as otherwise, e.g. for Pt/Co/Alo_x devices, the SOTs will be overestimated.

The detailed calculations are given in Appendix A.3.
3.3. Harmonic Hall voltage measurements

Spin-orbit torque analysis

In Fig. 3.7d-f we show the second harmonic signals obtained simultaneously to the first harmonic signals discussed in the previous section. For the analysis of the second harmonic signal we make use of the previously determined $R_{AHE}$, $R_{PHE}$, and $\theta(B_{ext})$ and we remind that $\varphi = \varphi_B$ for samples with uniaxial PMA. Different methods were developed for the evaluation of the effective SOT fields, $B_{FL}$ and $B_{DL}$ (Refs. 37, 43, 174). In the following, we will focus on two different methods which were developed in the course of this work.

One of them is to fit the measured second harmonic signal by Eq. (3.9). This method has the great advantage that it is possible to access $B_{FL}$, $B_{DL}$ as well as $R_{2\omega}$ from a single measurement if $\varphi_B \neq 0^\circ, 45^\circ, 90^\circ$ (see Appendix C.1 for the MATLAB code), which is very convenient for fast characterization. However, it is crucial to calibrate $\varphi_B$ precisely as all the $\varphi$-terms in Eq. (3.9) depend strongly on the accuracy of the calibration and may lead to wrong results otherwise. In order to improve the robustness of the fit, it is possible to fit measurements with different $\varphi_B$ simultaneously. Obviously, the big drawback is that the result can only be as good as the model which describes the underlying physics. In some cases, higher order terms (compare Eqs. (A.3), (A.4)) may have to be taken into account in order to get a satisfactory result. Figure 3.9a shows the measured data and the corresponding fit of the second harmonic curve. The analysis is based on the data shown in Figs. 3.7d-f, which for simplicity, were analysed for positive fields only. From the fit we find $B_{DL} = 18$ mT and $B_{FL} = B_{FL0} + B_{FL2}\sin^2(\theta)$ with $B_{FL0} = 8$ mT and $B_{FL2} = 6.5$ mT per $10^8$ A cm$^{-2}$.

The second method is based on the solution of a matrix equation which reveals the SOT amplitudes as a function of $B_{ext}$. In the simplest case, this is done by conducting a field scan measurement at $\varphi_B = 45^\circ$ and $\varphi_B = 315^\circ$, respectively\textsuperscript{IV}. At these angles the second term of Eq. (3.9) vanishes and hence, Eq. (3.9) can be simplified using the identity

$$\frac{dR^\omega_H}{dB_{ext}} = [R_{AHE} - 2R_{PHE}\cos(\theta)\sin(2\varphi)] \frac{d\cos(\theta)}{dB_{ext}}, \quad (3.12)$$

to

$$R_{H}^{2\omega} = \frac{dR_H^\omega}{dB_{ext} \sin(\theta_B - \theta)} \left[ B_{FL}\cos(\varphi) + B_{DL}\cos(\varphi) \right] + R_{\nabla T}. \quad (3.13)$$

\textsuperscript{IV}Using symmetry arguments it is possible to reconstruct the $\varphi_B = 315^\circ$ measurement from the measurement conducted at $\varphi_B = 45^\circ$ which reduces the number of measurements.
3.3. Harmonic Hall voltage measurements

Figure 3.9 | Analysis of second harmonic signal. a, Symbols show the raw $R_{HH}^{2\omega}$ signal obtained from field scans with $\theta_B = 86^\circ$ and $\varphi_B = 0^\circ, 45^\circ, 90^\circ$ (see Fig. 3.7d-f) while the result of a fit with Eq. (3.9) is illustrated by the corresponding solid lines. b, Result of the SOT analysis (red: $B^{FL}$, black: $B^{DL}$) based on the fitting procedure (solid lines) and the matrix method (squares and circles). The resulting SOTs obtained in the low field limit (stars) and the high field limit (pentagon) are shown as a comparison. From the fitting procedure we find $B^{DL} = 18 \text{ mT}$ and $B^{FL} = B^{FL0} + B^{FL2} \sin^2(\theta)$ with $B^{FL0} = 8 \text{ mT}$ and $B^{FL2} = 6.5 \text{ mT per } 10^8 \text{ A cm}^{-2}$ (compare Eq. (A.3)).

The system of equation, which has to be solved for $B^{FL(DL)}$ is

$$\left( \begin{array}{c} R_{H,45}^{2\omega} \\ R_{H,315}^{2\omega} \end{array} \right) = \frac{1}{2} \left( \begin{array}{cc} A_{45} & B_{45} \\ A_{315} & B_{315} \end{array} \right) \left( \begin{array}{c} B^{FL} \\ B^{DL} \end{array} \right) + \left( \begin{array}{c} R_{HT,45}^{2\omega} \\ R_{HT,315}^{2\omega} \end{array} \right).$$  \hspace{1cm} (3.14)

Here, $R_{H,i}^{2\omega}$ are the measured second harmonic signals, $R_{HT,i}^{2\omega}$ are the thermal contributions to the second harmonics according to Eq. (3.11) and $i = 45^\circ$ and $315^\circ$. From a comparison with Eq. (3.9) we find the coefficients

$$A_i = \frac{dR_{H,i}^{2\omega}}{dB_{ext} \sin (\theta_B - \theta)} \frac{1}{\cos (\theta) \sin (\varphi_i)},$$  \hspace{1cm} (3.15)

and

$$B_i = \frac{dR_{H,i}^{2\omega}}{dB_{ext} \sin (\theta_B - \theta)} \frac{1}{\cos (\varphi_i)}.$$  \hspace{1cm} (3.16)

One has to consider that only $B^{i}_{\theta}$ is evaluated, as the second term in Eq. (3.9) vanishes and $B^{i}_{\varphi}$ cannot be estimated\footnote{If we conduct four independent measurements ($\varphi_B = 0^\circ, 45^\circ, 90^\circ, 315^\circ$) Eq. 3.14 has to be expanded and we can access both $B^{i}_{\theta}$ and $B^{i}_{\varphi}$.}. Contrary to the analysis using the fitting procedure, $R_{HT}^{2\omega}$ has to be determined in a separate measurement. This can be done by exploiting the different field dependence of $R_{HT}^{2\omega}$, which does not depend
3.4 Scanning transmission x-ray microscopy

Table 3.2 | Typical parameters for a Pt(5 nm)/Co(1 nm)/AlO$_x$ sample deposited on a Si/Si$_3$N$_4$ substrate according to Tab. 3.1. $\xi^i_{DL}$ was calculated according to Eq. (2.12). *For a current density of $10^8$ A cm$^{-2}$.

<table>
<thead>
<tr>
<th>$R_{AH}$</th>
<th>$R_{PH}$</th>
<th>$B_k$</th>
<th>$B_{DL}$</th>
<th>$B^E_0$</th>
<th>$B^E_2$</th>
<th>$\rho$</th>
<th>$M_s$</th>
<th>$\xi^i_{DL}$</th>
</tr>
</thead>
<tbody>
<tr>
<td>(\Omega)</td>
<td>(\Omega)</td>
<td>(mT)</td>
<td>(mT*)</td>
<td>(mT*)</td>
<td>(\mu\Omega cm)</td>
<td>(kA m$^{-2}$)</td>
<td>()</td>
<td></td>
</tr>
<tr>
<td>0.26</td>
<td>0.09</td>
<td>680</td>
<td>18</td>
<td>8</td>
<td>6.5</td>
<td>25</td>
<td>925</td>
<td>0.04</td>
</tr>
</tbody>
</table>

The advantage of this method, however, is that dependences of $B^I_{\theta,\phi}$ on $\theta$ appear naturally in the solution of Eq. (3.14). Figure 3.9b shows $B^I_{\theta}$, the solution of Eq. (3.14), as a function of $\sin^2(\theta)$. The red circles correspond to $B^F_0$ and the black squares to $B^F_2$. Additionally, we indicate the result of the SOT analysis obtained in the low field limit$^{103}$ by the red and black stars and the resulting $B^F_0$ in the high field limit$^{174}$ by the black pentagon.

Table 3.2 summarizes the properties of the prototype Pt(5 nm)/Co(1 nm)/AlO$_x$ samples, grown according to Section 3.1.1 and fabricated into a 5 $\mu$m wide Hall bar (Fig. 3.2). We made sure that the deposition as well as the fabrication process is reproducible and that similar results are obtained from batch to batch.

Appendix C.2 shows a typical MATLAB script for the SOT-analysis based on the matrix method with four independent field scan measurements conducted at $\varphi_B = 0^\circ, 45^\circ, 90^\circ$ and $315^\circ$, hence, both $B^I_{\theta}$ and $B^I_{\phi}$ are accessible. See Appendix A.4 for the detailed calculation of the matrix method with four measurements.

3.4 Scanning transmission x-ray microscopy

In order to conduct time-resolved SOT-induced magnetization switching experiments we used the scanning transmission x-ray microscope (STXM) at the Pol-Lux beam line of the Swiss Light Source$^{175}$ implementing a current pump/x-ray probe measurement scheme. Figure 3.10 shows the schematic working principle of the STXM. A Fresnel zone plate focuses the monochromatic x-ray beam through an order selecting apperture (OSA) onto the sample, which is scanned through the x-ray focus while an avalanche photodiode collects the transmitted photon intensity. The magnetization state is probed by exploiting the x-ray magnetic circular dichroism (XMCD)$^{176,177}$ at the Co L$_3$ absorption edge ($E = 779$ eV) using left circularly polarized light at normal incidence.
Figure 3.10 | Schematic of the experimental setup. a, A circularly polarized x-ray beam is focused to a 25 nm spot by a zone plate and an order selecting aperture (OSA). The transmission of the x-rays through the sample is monitored by an avalanche photodiode as the sample stage is scanned in steps of 12-25 nm across the x-ray focus. b, Scheme of the pump current circuit. c, Picture of a sample mounted on the impedance matched wave guide. Below the sample the 50 Ω SMD-resistor connected in parallel with the sample can be seen.

The XMCD contrast depends on the relative orientation of the x-ray polarization vector and the magnetization component along the polarization vector. The incident circular polarized x-ray photons are absorbed by core electrons (e.g. \( p_{3/2} \) level, L\(_3\) absorption edge) and transfer their angular momentum through the SOC to the generated photoelectrons. As the angular momentum of left- and right-circular polarized photons is opposite, the emitted photoelectrons will have an opposite spin in the two cases. The spin-polarized photoelectrons are absorbed in empty states of the spin-split valence d shell. Given that spin-flips of the photoelectrons are forbidden, a spin-up (down) photoelectron can only be excited into an empty spin-up (down) d state. Hence the x-ray absorption cross section is proportional to the imbalance of the spin-up and spin-down electrons of the exchange split d band and, therefore, to the magnetization direction with respect to the x-ray polarization vector. As we work under normal inci-
3.4. Scanning transmission x-ray microscopy

dence, the difference in x-ray absorption for left and right circularly polarized x-rays, that is the XMCD, is directly related to the out-of-plane component of the magnetization ($M_z$) of the Co dots.

Figure 3.11a shows a schematic representation of the samples used for the STXM experiments. All the STXM experiments described herein were performed on Pt(5 nm)/Co(1 nm)/AlO$_x$ stacks which were deposited by dc magnetron sputtering according to Section 3.1.1. Due to the nature of the experiment, the Pt/Co/AlO$_x$ layers have been grown on transparent Si$_3$N$_4$ membranes. The membrane windows have a thickness of 200 nm. Following the description in Section 3.1.3 we fabricated 500 nm large Co dots on top of 750 nm wide Pt lines contacted by Au leads. Figure 3.11b shows a static image of a Co dot and the Pt line obtained by scanning the sample under the focused x-ray beam. The contrast provided by the XMCD allows to distinguish between the two out-of-plane remanent magnetization states, as shown in Fig. 3.11c. Typical dot images are obtained by recording the x-ray intensity over a two-dimensional $64 \times 64$ pixel array with about 25 nm lateral resolution.

Time-resolved images of the Co magnetization were recorded stroboscopically employing the time structure of the synchrotron radiation, which consists of 70 ps wide x-ray pulses with a repetition rate of 500 MHz (see Fig. 3.12a). Simultaneously the sample is pumped with a bipolar current pulse sequence that has a total length of 102.1 ns. Each unit-sequence includes a positive and negative pulse spaced by 50 ns, the length and amplitude of which can be tuned independently from each other, as indicated in Fig. 3.12b. The amplitudes of the two current pulses are adjusted in order to trigger SOT-induced magnetization switching\textsuperscript{VI}. Every 2 ns the magnetization is probed by the x-ray pulses (Fig. 3.12c) which results in a time resolution of 100 ps after the injection of 20 successive unit-sequences by routing the intensity of each probing event to a counting register (Fig. 3.12d). The current pulse sequences run continuously during the measurements, synchronized to the x-ray pulses, and are monitored by means of a -20 dB pick-off tee connected to an oscilloscope. The pump current circuit is depicted in Fig. 3.10b. Pulse reflections are minimized by mounting the sample on an impedance matched wave guide and connecting it in parallel to a 50 $\Omega$ SMD-resistor (Fig. 3.10c). Additionally, the STXM setup allows us to apply an in-plane magnetic field $B_x$ along the current direction in order to uniquely define the SOT-induced switching polarity\textsuperscript{20}.

\textsuperscript{VI}This approach allows us to use a similar set/reset pulse approach as for the all-electrical switching experiments described in Section 3.2.2.
3.5. Micromagnetic simulations

In the following section we introduce the micromagnetics simulations conducted as part of this thesis. All the simulations are based on the object oriented micromagnetic framework (OOMMF) code including the DMI extension module as well as an SOT module written by myself. During micromagnetic simulations the LLG-equation (see Eq. (2.16)) is solved for magnetic moments evenly spaced on a rectangular grid with a unit cell size of 4 nm × 4 nm × 1 nm.

The magnetic moment interact with each other due to the exchange coupling, the dipolar field and the DMI. (see Section 2.1). If not stated differently, we use the following material parameters for all micromagnetic simu-
3.5. Micromagnetic simulations

**Figure 3.12 | Pump probe measurement scheme.** a, Illustration of the x-ray time structure of the synchrotron (70 ps wide x-ray pulses with repetition rate of 500 MHz). b, The sample is pumped by means of a bipolar current pulse sequence with amplitudes optimized for reproducible SOT-induced magnetization switching. Each sequence is 102.1 ns long and consists of two pulses separated by 50 ns. c, The magnetization state is probed every 2 ns by an x-ray pulse, indicated by the black crosses. The transmitted x-ray intensity of each event is measured by the avalanche photodiode and stored in a counting register. d, After injection of 20 unit-sequences a time resolution of 100 ps is obtain by resorting the measured intensities in the counting register.

Damping: saturation magnetization $M_s = 900 \text{kA m}^{-1}$, exchange coupling constant $A_{ex} = 10^{-11} \text{J m}^{-1}$, uniaxial anisotropy energy $K_u = 657 \text{kJ m}^{-3}$, DMI $D = 1.2 \text{mJ m}^{-2}$, and damping constant $\alpha = 0.5$. The magnitudes of the damping- and field-like SOTs correspond to the values measured by the harmonic Hall voltage method given in Table 3.2. Typical current (pulse) amplitudes are $j_p = 4.5 \cdot 10^8 \text{A cm}^{-2}$ and a bias field of $B_x = 93 \text{mT}$ is applied if not stated differently. All simulations were carried out at zero temperature and the sample thickness was in all cases 1 nm. Depending on the simulations, different shapes (and initial states) were used. The results presented in Chapter 4.3 were conducted on 500 nm wide circular dots starting from homogeneous initial states. The simulations shown in Chapter 5 are based on different initial states and shapes, which are defined by a bitmap image. Using colour coding of such a bitmap image, material parameters, initial magnetization states and current densities can be defined for each simulation cell separately.
3.5. Micromagnetic simulations

The OOMMF input file used for the simulations described in Section 5 is reported in Appendix E.
4 Spatially and time-resolved magnetization dynamics

This chapter is based on the publication “Spatially and time-resolved magnetization dynamics driven by spin-orbit torque” (Ref. 79). I contributed to the planning of the experiment, fabricated the samples, implemented and performed the x-ray measurements (together with many other colleagues) and analysed the data. Furthermore, I conducted the electrical measurements on replica dots and the micromagnetic simulations and was one of the main contributors to the manuscript.

4.1 All-electrical measurements of magnetization switching

In the following we show the results obtained during all-electrical magnetization switching experiments conducted on Pt/Co/AlO$_x$ samples patterned into 500 nm wide Co dots. This type of samples are replicas of those employed for the STXM measurements, which we used to characterize the magnetic properties by probing the anomalous Hall effect.

4.1.1 Spin-orbit torque induced switching of circular dots

Following the procedure given in Section 3.2.2 we conduct all-electrical pulsed magnetization switching experiments on the 500nm wide Co dots. Figure 4.1a shows an atomic force microscope image of such a Co dot patterned on top of a Pt Hall cross deposited on a 200 nm thick Si$_3$N$_4$ membrane window (according to Sections 3.1.1 and 3.1.3) intended for the STXM-measurements discussed in Section 4.2. The Co dot is 1 nm thick, capped by 2 nm of oxidized Al, and the Pt line is 5 nm thick. It is important to note that the results obtained during these all-electrical switching measurements do not depend on the Si$_3$N$_4$ membrane windows, but could be reproduced on normal Si/SiO$_2$-chips.

Figure 4.1b shows the Hall resistance $R_H$ measured as a function of the external field $B_{ext}$ applied out-of-plane ($\theta_B = 0^\circ$) and nearly in-plane ($\theta_B = 84^\circ$). The square loop measured at $\theta_B = 0^\circ$ shows that the dots have strong PMA. With Eq. (3.5) we estimate the effective anisotropy field $B_k \approx 660$ mT from the $\theta_B = 84^\circ$ measurement.

Figure 4.1c shows pulsed switching measurements as a function of the pulse
4.1. All-electrical measurements of magnetization switching

4.1.2 Role of the field-like torque

The FL torque is normally neglected in models of SOT-induced magnetization switching, however, we show in the following that the reversal process is strongly influenced by the FL torque. We discussed in Section 2.3 that SOT-induced switching is believed to be most efficient if the bias-field is applied along the current line, this is for $\varphi_B = 0^\circ$ and $\theta_B = 90^\circ$. As $B^{FL}$ is parallel to the $y$-direction one could expect the switching to be more efficient if we compensate for it with the external field in order to obtain a pure $B_x$ field. Following this
4.2. Spatially and time-resolved measurements

Employing the STXM setup described in Section 3.4 we studied time-resolved SOT-induced magnetization dynamics in Co islands with perpendicular magnetization.

Figure 4.3a shows the XMCD time trace obtained by integrating the trans-

Figure 4.2 | Effect of an external field balancing the FL torque. a, Magnetization change measured by recording the difference of the Hall resistance before and after the pulse as a function of the applied voltage $U_p$ and the in-plane external field $B_y$. $\Delta R_H$ is proportional to the fraction of the dot area that has reversed its magnetization, averaged over 200 switching events. The $x$-component of the external field is fixed to 92 mT. b, Threshold voltage at 50% switching as a function of $B_y$. The vertical error bars represent the standard deviation of the threshold voltage over repeated switching trials. The horizontal error bars represent the maximum uncertainty of $B_y$ due to sample misalignment. Inset: schematic of the field direction. For a negative current, $B_y$ opposes (favours) $B^{FL}$ for $B_y > 0$ ($B_y < 0$).
4.2. Spatially and time-resolved measurements

mitted x-ray intensity over the entire area of the Co dot during the bipolar current pulse sequence (see Fig. 3.12), which was separately recorded by a fast oscilloscope in parallel with the sample (see Fig. 3.10) for the electrical circuit. The pulses in this sequence are 2 ns long, with a rise time of about 150 ps, and have a separation of 50 ns. The time trace shows that the magnetization switches from down to up during the first (negative) pulse and from up to down during the second (positive) pulse. The timescale over which \( M_z \) changes between two states with opposite saturation magnetization coincides with the duration of the current pulses, indicating that the magnetization is fully reversed within the

![Figure 4.3](image_url)

**Figure 4.3 | Time-resolved magnetization switching.** XMCD time traces showing reversal of the magnetization during the injection of 2 ns long current pulses of opposite polarity at a, \( B_x = 94 \text{ mT} \) and b, \( B_x = -124 \text{ mT} \). The sign of the XMCD signal is positive for \( M_z < 0 \) and negative for \( M_z > 0 \). c, Asymmetric set-reset sequence in which the first pulse is 2 ns long and has amplitude \( U_p = -3.7 \text{ V} \) while the second pulse is 0.8 ns long and has amplitude +4.5 V. d, Fast toggle sequence consisting of two 1 ns long current pulses of opposite polarity separated by 1 ns with in-plane field \( B_x = 92 \text{ mT} \). \( U_p = 1 \text{ V} \) corresponds to a current density \( j_p = 8.4 \cdot 10^7 \text{ A cm}^{-2} \) in the Pt line. Each time trace is the average of about \( 1.2 \cdot 10^{11} \) consecutive switching cycles recorded with a 20 MHz repetition rate over 100 minutes.
pulse interval, with no significant delay or after-pulse relaxation. The response of the magnetization to the current pulses inverts upon changing the sign of $B_x$, as shown in Fig. 4.3b, consistent with previous reports of SOT-induced switching of perpendicularly magnetized layers. Further, we find that the switching speed increases by increasing the current amplitude, allowing for full magnetization reversal within a sub-ns current pulse (Fig. 4.3c). This speed, and the absence of after-effects, enable the realization of very fast magnetic-writing cycles. One such cycle is shown in Fig. 4.3d, where we toggle the magnetization between up and down states using a 1 ns on | 1 ns off | 1 ns on pulse sequence. Remarkably, our samples survive uninterrupted pulse sequences for hours without appreciable magnetic or electrical degradation. We tested more than $10^{12}$ consecutive and successful switching events, at current densities of the order of $2 - 4 \cdot 10^8$ A cm$^{-2}$. Thus, the combination of speed and endurance revealed by our measurements is extremely promising for the operation of SOT based memory and logic devices.

The pulse voltages and width used for reliable switching are comparable to those used for the all-electrical switching measurements on the replica dots (see Section 4.1.1). The main difference being that the $U_p$ values in the Hall measurements are about 20% larger compared to the x-ray measurements due to the dispersion of the current in the Hall cross.

4.2.1 Spatial evolution of the magnetization during current injection

We next focus on the transient magnetic configurations and the mechanisms leading to magnetization reversal. Figure 4.4a shows four series of consecutive images recorded at time intervals of 100 ps during the switching process, corresponding to the four possible combinations of current and field polarity. The magnetization reverses by domain nucleation and propagation in all cases, with no appreciable incubation delay. Although the stroboscopic character of our measurements does not allow us to investigate stochastic effects, the observation of a clear DW front moving from a fixed nucleation point on one side of the sample to the opposite side, as shown by the red dots and green arrows in Fig. 4.4a, indicates that the reversal process is reproducible and deterministic. Furthermore, as we argue in the following, such a reversal scheme is unique to SOTs. Our data show that domain nucleation takes place at the edge of the sample where the DMI and $B_x$ concur to tilt the magnetization towards the current direction, thereby confirming the prediction of recent micromagnetic models (see Section 2.3.2 for more details about the micromagnetic models).
4.2. Spatially and time-resolved measurements

Figure 4.4 | Evolution of the magnetization during the switching process. 

a, Images taken at intervals of 100 ps during the injection of 2 ns long current pulses. Rows (I, II) and (III, IV) correspond to the time traces shown in Fig. 4.3a and b, respectively. The red dots indicate the DW nucleation point and the green arrows its propagation direction. The images are low-pass filtered for better contrast (see Fig. B.1 and the Appendix B for the raw data). 

b, Schematic representation of the observed DW nucleation and propagation geometry. Illustration of the nucleation process corresponding to case (II).

c, Canting of the magnetization at the dot edges induced by the DMI.

d, Breaking of the canting symmetry induced by $B_{x}$. 

e, Action of $B^{DL}_{F}$ and $B^{DL}_{L}$.

However, the concerted action of $B^{DL}$, DMI, and $B_{x}$ only leads to a left/right asymmetry, similar to the asymmetric nucleation induced by a perpendicular magnetic field, whereas we observe that the domain nucleation site alternates between the four quadrants shown in Fig. 4.4b, with an additional top/bottom asymmetry.

4.2.2 Edge nucleation

To explain this additional asymmetry, we have to analyze the effects of the static and dynamic fields in more detail. Pt/Co/AlO$_x$ layers have positive DMI, which stabilizes left-handed Néel DWs and induces canting of the magnetization at the edge of the dot. In the absence of current and external field, the magnetic
moments are symmetrically canted inwards (outwards) of the dot for $M_z > 0$ ($M_z < 0$), as illustrated in Fig. 4.4c. Upon applying a static field $B_x$, the canting angle increases on one side, while it decreases on the other (Fig. 4.4d), favouring domain nucleation on the side where the canting is larger. The injection of a positive current pulse generates an effective DL torque with a polar component $B_{DL}^\theta$ that points toward the current direction (purple arrow in Fig. 4.4e). For a positive $B_x$, $B_{DL}^\theta$ thus leads to left nucleation if $M_z > 0$ and no nucleation if $M_z < 0$. A similar reasoning is valid for all four combinations of field and current polarity leading to magnetization switching, which explains the left/right asymmetry observed in our data. The top/bottom asymmetry, however, can be explained only if an additional torque plays a significant role in the nucleation process. In the following, we argue that the FL torque, combined with the canting induced at the sample edges by the DMI, accounts for such an asymmetry. Because the corresponding effective field $B^{FL}$ points along the $y$-direction and has no projection along the easy axis, the effects of this torque are usually neglected in models of the SOT-induced DW dynamics. However, the polar component $B^{FL}_\theta \propto \cos \theta \sin \varphi$ points upward or downward depending on the sign of $M_y$, as illustrated by the blue arrows in Fig. 4.4e. Therefore, the rotation of the magnetization and the nucleation of a reverse domain are favoured whenever $B^{FL}_\theta$ and $B_{DL}^\theta$ are parallel to each other, as indicated by the red dots in Fig. 4.4, and hindered when they are antiparallel. This qualitative argument is supported by harmonic Hall voltage measurements of the FL and DL torques in our samples, which show that $B^{FL}$ and $B^{DL}$ have comparable amplitudes of about 18 mT per $10^8 \text{ A cm}^{-2}$ (Table 3.2), and by both macrospin and micromagnetic simulations that take these torques into account. We note that, in principle, also the $z$-component of the Oersted field generated by the current flowing in the Pt line can induce a similar top/bottom asymmetry as reported here and assist the nucleation process. However, while Oersted field-assisted switching is of interest for device applications, we find that its effects are minor in the present case, as the closest edge of the Co dot is about 125 nm from the edge of the current line and the Oersted field is significantly smaller compared to the FL torque (Fig. 4.12).

4.2.3 Macrospin model of magnetization reversal at different edge positions

We provide a simplified and intuitive model of the edge nucleation by simulating the reversal process of four non-interacting single spins located at positions (1) - (4) around the Co dot, as illustrated in Fig. 4.5a. The DMI is included as an
4.2. Spatially and time-resolved measurements

effective field by considering the DMI-energy

\[ E_{DMI} = \mathbf{D} \cdot (\mathbf{M}_i \times \mathbf{S}), \quad (4.1) \]

where \( \mathbf{M}_i \) is the magnetization vector of the spin at position \( i \), \( \mathbf{S} \) represents a fictitious central spin of fixed orientation and \( \mathbf{D} \) is the DMI-coupling vector. \( \mathbf{D} \) lies in the sample plane and is perpendicular to the vector connecting the spins \( \mathbf{M}_i \) and \( \mathbf{S} \). It can thus be parametrized by an angle \( \varphi_i \) as \( \mathbf{D} = D(\sin \varphi_i, -\cos \varphi_i, 0) \).

We then make use of the general relation \( \mathbf{B}_{DMI} = -\nabla \mathbf{M} E_{DMI} \) to express the DMI-energy as an effective field

\[ \mathbf{B}_{DMI,i} = \text{sgn} (S_z) 2D \begin{pmatrix} \cos \varphi_i \\ \sin \varphi_i \\ 0 \end{pmatrix}. \quad (4.2) \]

Finally, we perform a time-integral of the LLG equation (compare Section 2.3.1) to separately model the evolution of the magnetization at points (1) - (4) under the action of \( B_{DMI,i} \), \( B_x = 100 \text{ mT} \), effective anisotropy field \( B_k = 657 \text{ mT} \), and the DL and FL torques given in Table 3.2. We choose the DMI-coupling strength \( D \) such that it induces a spin canting at the sample edge similar to that reported in Ref. 65. The canting angle at positions (1) and (2) under a positive bias field of \( B_x = 100 \text{ mT} \) is \( \approx 18^\circ \) for \( D = 60 \text{ mT} \), whereas it is \( \approx 8^\circ \) at positions (3) and (4). For brevity, we limit the discussion to \( B_x > 0 \) and positive current pulses, as indicated in Fig. 4.5a.

We consider first the case \( S_z > 0 \), which represents the dot in the up-state. For a given current, the simulated time traces of the magnetization at positions (1) - (4), reported in Fig. 4.5b, show that only the spin located at position (1) reverses, whereas, by inverting the sign of \( B^{FL} \) and keeping \( B^{DL} \) constant, only the spin at position (2) reverses. This behavior is in agreement with the nucleation point observed experimentally (Fig. 4.4b (II)). Inverting the sign of the central spin, such that \( S_z < 0 \), changes the direction of the effective DMI-field and makes position (3) equivalent to position (1) in the former case (\( S_z > 0 \)). The effect of the FL torque at different positions can be simply rationalized by considering the \( \theta \)-component of \( B^{FL} \):

\[ B^{FL}_\theta = B^{FL} \cdot \mathbf{e}_\theta = \begin{pmatrix} 0 \\ B^{FL} \\ 0 \end{pmatrix} \cdot \begin{pmatrix} \cos \theta \cos \varphi_i \\ \cos \theta \sin \varphi_i \\ -\sin \theta \end{pmatrix} = B^{FL} \cos \theta \sin \varphi_i, \quad (4.3) \]
4.2. Spatially and time-resolved measurements

Figure 4.5 | Macrospin model of edge magnetization reversal. 

**a**, Schematic representation of four non-interacting spins located at distinct positions around the dot. The DMI is taken into account as an effective in-plane field, the sign of which depends on the orientation of the fictitious central spin. **b**, z-component of the four macrospins during the application of a current pulse for $B^{FL} > 0$ (upper panel, corresponding to the experimental situation) and $B^{FL} < 0$ (lower panel).

Figure 4.6 | Effect of the external in-plane field on the switching speed. 

**a**, Measured XMCD time traces as a function of in-plane field $B_x$. **b**, Average slope of the rising and falling edges during the reversal plotted as function of $B_x$.

where we see that $B_\theta^{FL}$ depends on $\varphi_i$ and points towards the $xy$-plane at position (1) and in the opposite direction at position (2), thus favoring the reversal of spin (1) relative to (2).
4.2. Spatially and time-resolved measurements

4.2.4 Dependence of the switching speed on the external field

Figure 4.6a shows the XMCD time traces of a Co dot during the injection of 2 ns long current pulses of amplitudes $U_p = -3.35 \text{ V}$ and $3.42 \text{ V}$ at different values of the applied field $B_x$. In the following, we will limit the discussion to negative fields but emphasize that the same arguments are consistent with positive bias fields. We find that $B_x$ has an influence on the reversal speed, namely that switching occurs faster at high field values. This is shown in Fig. 4.6b where the average slope of the XMCD time traces during the reversal is plotted against $B_x$. With increasing $B_x$, $M_z$ decreases while $|M_x|$ increases on one side of the dot, which favours magnetization reversal. Our findings are supported by micromagnetic simulations (see Section 4.3.2) in which we analyze the influence of $B_x$ on the reversal speed.

4.2.5 Partial switching at sub-threshold current amplitude

By calculating the time required for the DW to cover the central region of the dots, we estimate that the DW velocity is of the order of $400 \text{ m s}^{-1}$, corresponding to about $100 \text{ m s}^{-1}$ per $10^8 \text{ A cm}^{-2}$ of injected current, in agreement with quasi-static measurements of DW displacements. As magnetization reversal is deterministic and achieved by a single DW traversing the entire sample, the switching timescale is expected to be directly proportional to the lateral sample size. This can lead to switching times of less than 200 ps in structures that are smaller than 100 nm. Moreover, we find that pulses that are either shorter or weaker in amplitude compared to the threshold values required to achieve full switching consistently lead to the reversal of a fraction of the dot area. Figure 4.7a shows the result of a series of switching measurements taken at increasing values of the voltage applied to the current line. Each frame is a differential image showing the average dot area that reversibly switches the magnetization upon applying positive and negative pulses. The reversed dot area increases monotonically with the pulse amplitude, as illustrated in Fig. 4.7b, and correlates well with the remanent magnetization measured by the anomalous Hall effect on a replica dot (compare Section 4.1.1). These measurements show that the critical switching current is mostly dependent on the DW mobility and sample dimensions rather than on the initial nucleation barrier. Further, our results agree with the absence of DW inertia reported in Pt/Co layers, consistently with the fast damping of magnetic excitations in SOT devices, and show that partial but reliable switching can be obtained also when working below the
4.2. Spatially and time-resolved measurements

Figure 4.7 | Partial switching induced by sub-threshold current pulses. a, Differential images showing the extent of magnetization reversal (in white) for pulses of increasing voltage amplitude. The differential contrast is obtained by averaging all the frames in a time-sequence after positive pulse injection and subtracting the average of all the frames after negative pulse injection. b, Comparison between the fractional reversed area estimated from the images in a, and the all-electrical switching measured by the Hall resistance on a replica dot (compared Section 4.1.1). The Hall resistance data are averaged over 200 pulse cycles. The different voltage scale for the two measurements is due to the dispersion of the current in the branches of the Hall cross, which is absent in the x-ray measurements.

Figure 4.8 | Switching in the presence of defects. Images taken at intervals of 100 ps during the injection of 2 ns long current pulses. DW nucleation is observed at different sites. The point at which the reversed domain front collapses (orange dot) is off-centred, in agreement with the preferential DW propagation direction discussed in the main text (green arrows).

current amplitude required for full switching.

4.2.6 Switching of defective dots

All of the measured dots show reproducible switching behavior rather than thermally-induced stochastic reversal, which would result in uniform STXM contrast during the reversal process. However, approximately 50% of the dots
4.3. Micromagnetic simulations of circular dot switching

showed defect-promoted switching, likely due to inhomogeneities occurred during the patterning process. Such samples are distinguished from the ones presented in Fig. 4.4 by the fact that nucleation always starts from a given point or region rather than alternating between the four dot quadrants. Figure 4.8 shows a case of defect-promoted switching in which domain nucleation starts at several points around the Co dot edge. Due to the stroboscopic nature of our measurements, we cannot tell whether the nucleation starts simultaneously at different locations or is randomly distributed over the dot edge. However, we do observe that once a domain has nucleated, the favoured DW propagation direction remains the same (green arrows in Figure 4.8) and is comparable to the switching behaviour reported in Fig. 4.4 of the main text. The point at which the reversed domain front collapses (orange dot) is therefore off-centred. Importantly, we find that in all cases the reversal process is robust with respect to the presence of defects, which makes SOT-induced magnetization switching very versatile for applications.

4.3 Micromagnetic simulations of circular dot switching

In the following we present the results obtained with micromagnetic simulations based on the integration on the LLG-equation as explained in Section 3.5.

4.3.1 Dynamic domain wall propagation

A striking feature observed in Fig. 4.4 is that the propagating domain front is tilted relative to the current direction, with a tilt angle of about 45° that changes in steps of 90° depending on the up/down or down/up DW configuration and the sign of the current. According to recent studies, current-induced DW tilting is a telltale signature of the DMI in perpendicular magnetized nanotacks (see also Section 2.3.2 and 5). However, the tilt angle in Fig. 4.4 is opposite to that predicted by micromagnetic models of Pt/Co heterostructures and observed by MOKE microscopy in Pt/Co/Ni/Co racetracks. More specifically, the angle between the DW normal and the current direction is ≈−45° for a left-handed up-down wall (↑←↓) at positive current (see panel IV in Fig. 4.4b) rather than ≈+45° as reported in previous studies. We believe that this inconsistency stems from the neglect of the FL torque in the micromagnetic models of current-induced DW motion as well as from the time-resolved nature of our measurements (see Chapter 5). The tilt angle in the Pt/Co/AIOx dots is in fact analogous to that induced by an external in-plane magnetic field $B_y$, which leads to a rotation of the internal DW magnetization away from the x-axis in order to
4.3. Micromagnetic simulations of circular dot switching

recover the Néel configuration favoured by the DMI\textsuperscript{157,166}. Our micromagnetic simulations, which include $B^\text{FL}$ as well as $B^\text{DL}$ and the DMI, correctly reproduce the observed dynamic tilt during DW propagation, as shown in Fig. 4.9. Moreover, the simulations indicate that the FL torque promotes faster DW propagation in the direction indicated by the arrow in Fig. 4.9b, which coincides with case II reported in Fig. 4.4 and may also explain the strong anisotropy of the DW velocity recently reported in extended Pt/Co/AlO\textsubscript{x} layers\textsuperscript{183}. We note further that the DW propagation direction is not related to the bias field $B_x$, the main purpose of which is to break the spin-canting symmetry due to the DMI, whereas the DW velocity increases with $B_x$ (Figs. 4.6, 4.10). As the tilt angle depends on the $B^\text{FL}/B^\text{DL}$ ratio, the opposite tilt relative to the Pt/Co/Ni/Co racetrack\textsuperscript{162} may be explained by the different SOT amplitudes in this system. A more important difference, however, is that we probe the dynamic structure of the DW during current injection rather than after the current-induced displacement. Starting from a homogeneous magnetization state, we image the fastest domain front sweeping through the sample, which has opposite tilt with respect to the slowest front that survives in steady state conditions\textsuperscript{162} (see Fig. 5.3).
4.3. Micromagnetic simulations of circular dot switching

Accordingly, we find that the domain front in our measurements is orthogonal to the direction of largest DW velocity recently reported for Pt/Co/AlO$_x$\cite{185}. These findings show that magnetization switching is boosted in the Pt/Co/AlO$_x$ dots by the favourable combination of the domain nucleation symmetry and DW propagation direction, which is such that the fastest domain front can sweep unhindered across the full extension of the dot. As an example of a less favourable case, we simulated the effect of a negative $B^{FL}$ (Fig. 4.9a). Such a field would move the domain nucleation point to the opposite edge of the dot, while the direction of DW propagation along the $x$-axis would remain unaltered, leading to a slower reversal dynamics. We note that the presence of defects can also alter the DW dynamics, but does not prevent switching (Figs. 4.8, 4.13).

4.3.2 Dependence of the switching speed on the external field

Figure 4.10a shows the simulated time traces of $M_z$ following the injection of a current pulse of amplitude $j_p = 4.5 \cdot 10^8$ A cm$^{-2}$ for different values of $B_x$. The simulations show that there is a threshold bias field $63 \text{ mT} < B_x < 73 \text{ mT}$ below which reversal does not take place. Above this field, the DW velocity increases with $B_x$ as the magnetization tilts more towards the sample plane and the energy barrier for DW propagation reduces, in agreement with the experiment. We further note that the experimental threshold field can be significantly lower compared to the simulations owing to Joule heating, which assists the nucleation of reversed domains.

![Figure 4.10](image)

**Figure 4.10 | Effect of the external in-plane field on the switching speed.** a, Simulated time traces of $M_z$ as a function of $B_x$. The slow dynamics after the current pulse is due the DW reorientation after incomplete magnetization reversal. b, Reversal rate $\frac{dM_z}{dt}$ evaluated for the traces shown in a during the current pulse.
Switching speed as a function of current amplitude

Micromagnetic simulations performed as a function of $j_p$ show that complete reversal of the dot is achieved faster for larger current densities (Fig. 4.11a), in agreement with experimental observations. The simulations also evidence an initial tilting motion of the magnetization that precedes the nucleation of a reversed domain at the dot edge. This motion gives rise to a semi-precessional local minimum in the simulated time trace of $M_z$ (inset of Fig. 4.11b), which can be used to define an effective nucleation time $t^*_{N}$. This time decreases with increasing $j_p$, albeit weakly (Fig. 4.11b). In theory, such a $t^*_{N}$ can become a limiting factor for the switching speed of dots of very small size. In practice, however, we do not observe evidence of a distinct nucleation phase within the temporal resolution and sensitivity of our measurements. The question remains

Figure 4.11 | Simulations of magnetization reversal as a function of current amplitude and pulse length. a, Simulated time traces of $M_z$ as a function of $j_p$. b, Effective nucleation time $t^*_{N}$ defined by the duration of the initial tilt motion of the magnetization. Inset: Detail of the traces shown in a. The position of the asterisks defines $t^*_{N}$. c, Simulated time traces of $M_z$ as a function of pulses of length $\tau_p = 0.6$, 1.2 and 2 ns for $j_p = 4.5 \cdot 10^8$ A cm$^{-2}$. 
open if such a phase exists and can be experimentally resolved. Finally, we note that micromagnetic simulations performed as a function of pulse length at constant current density show that the switched dot area is roughly proportional to the pulse duration for pulses shorter than the time required to achieve full switching, as shown in Fig. 4.11c.

### 4.3.4 Effect of the Oersted field

The current flowing in the Pt line produces an Oersted field with y- and z-components given by

\[
B_{y}^O = -\frac{I}{wd} (U_y^+ - D_y^- + U_y^- - D_y^+) \cdot 10^{-3}, \tag{4.4}
\]

\[
B_{z}^O = -\frac{I}{wd} (U_z^+ - D_z^- - U_z^- + D_z^+) \cdot 10^{-3}, \tag{4.4}
\]

where \( I \) is the current, \( w \) is the width of the Pt line and \( d \) its thickness, respectively. The coefficients are defined as follows:

\[
U_y^\pm = 2(z + d) \arctan \left( \frac{w/2 \pm y}{z + d} \right) \left( \frac{w/2 \pm y}{z + d} \right)^2 + (z + d)^2, \tag{4.5}
\]

\[
D_y^\pm = 2z \arctan \left( \frac{w/2 \pm y}{z} \right) \left( \frac{w/2 \pm y}{z} \right)^2 + z^2, \tag{4.6}
\]

\[
U_z^\pm = 2(w/2 \pm y) \arctan \left( \frac{z + d}{w/2 \pm y} \right) \left( \frac{z + d}{w/2 \pm y} \right)^2 + (z + d)^2, \tag{4.7}
\]

\[
D_z^\pm = 2(w/2 \pm y) \arctan \left( \frac{z}{w/2 \pm y} \right) \left( \frac{z}{w/2 \pm y} \right)^2 + z^2, \tag{4.8}
\]

with \( z \) being the height above the surface of the Pt line.

Figure 4.12 shows the Oersted field calculated according to Eq. (4.4)-(4.8) for a current density of \( 1 \cdot 10^8 \text{A cm}^{-2} \) flowing through the Pt layer, at \( z = 0.5 \text{nm} \) above the Pt surface. The y-component \( B_{y}^O = -3 \text{ mT} \) is approximately constant over the dot surface and directed against \( B^{FL} \). The largest z-component is found at the two extrema of the dots closest to the edge of the Pt line, where
4.3. Micromagnetic simulations of circular dot switching

Figure 4.12 | Influence of the Oersted field on magnetization reversal. a, Calculated Oersted field according to Eq. (4.4) at a height of 0.5 nm above the Pt line for a positive current pulse of amplitude $1 \cdot 10^8 \text{ A/cm}^2$. b, Illustration of the nucleation process supported by the $z$-component of the Oersted field at positive and negative current. c, $\theta$-component of $B^{FL}$ and $B^{Oe}$ as a function of $\theta$ during a positive current pulse of amplitude $1 \cdot 10^8 \text{ A/cm}^2$ for a spin located at the position (1) and e, (2) of the dot (see diagram in Fig. 4.5). f, Simulated time traces for different combinations of $B^{Oe}$, $B^{FL}$ and $B^{DL}$. g, Corresponding snapshots during the magnetization reversal.

$|B^{Oe}| = 1.6 \text{ mT}$. The sign of $B^{Oe}$ is opposite on opposite sides of the dot, such that this field can also induce a top/bottom edge asymmetry, supporting the nucleation process induced by $B^{FL}$, as shown in Figs. 4.12b and c.
To analyze the role played by the Oersted field and the FL torque in the nucleation, we plot in Fig. 4.12d the $\theta$-components of $B^{FL}$ and $B^{Oe}$ as a function of $\theta$. The fields are evaluated at position (1) in Fig. 4.5, where the magnetization is tilted by the DMI and external field along $\varphi = 45^\circ$. At this position, the total $B_\theta$ is positive for all $\theta$, parallel to $e_\theta$ (Fig. 3.6), and thus favours the rotation of the magnetization from up to down. However, without $B^{FL}$ magnetization reversal would not be favoured at this position since $B^Oe_\theta$ is negative over a large range of $\theta$. The opposite argumentation is valid for the magnetization at position (2) in Fig. 4.5a, where the reversal is favoured by $B^Oe$ and hindered by $B^{FL}$ (Fig. 4.12e). As (1) corresponds to the experimentally observed nucleation, we conclude that $B^{FL}$ is the main cause of the nucleation asymmetry. It is important to note that $B^Oe_y$ is approximately constant over the dot area whereas $B^Oe_z$ increases rapidly near the Pt line edge. This implies that magnetization nucleation will be eventually also assisted by the Oersted field for Co dots with a diameter comparable to the Pt line width. To demonstrate the effect of $B^{Oe}$ on the switching process, we performed a series of micromagnetic simulations with

1) $B^{Oe} \neq 0$ and $B^{FL} \neq 0$,
2) $B^{Oe} = 0$ and $B^{FL} \neq 0$,
3) $B^Oe \neq 0$ and $B^{FL} = 0$
4) $B^Oe = 0$ and $B^{FL} = 0$, the results of which are reported in Figs. 4.12f and g. Without $B^{FL}$ the switching efficiency is strongly reduced, whereas the addition of $B^{Oe}$ to $B^{FL}$ has only a small influence on the reversal mechanism for our device geometry. Note also that the nucleation point moves up along the edge for $B^{FL} = 0$. Additional simulations show that the effect of the Oersted field becomes noticeable when the distance between the edges of the dot and of the Pt current line reduces to less than 50 nm.

### 4.3.5 Switching of defective dots

The robustness of SOT-switching relative to the presence of imperfections is confirmed also by micromagnetic simulations of defective dots. Figure 4.13 compares the switching of an ideal dot with uniform magnetic properties to that of dots with common defects such as regions with reduced magnetic anisotropy or saturation magnetization, as may arise due to the patterning process or thickness fluctuations. For test purposes, some defects are placed on the side of the dot where they would alter the nucleation mechanism described in Section 4.2.3.

For such localized edge defects, we observe that a lower magnetization hinders nucleation at the “right” edge of the dot, whereas a lower magnetic anisotropy favours nucleation at the “wrong” edge of the dot, as shown in Fig. 4.13b and c. In either case, however, we find that nucleation and DW propagation from the
4.3. Micromagnetic simulations of circular dot switching

**Figure 4.13 | Micromagnetic simulations of switching in the presence of defects.**

- **a**, Snapshots of the magnetic configuration at intervals of 100 ps in the case of a faultless dot. The saturation magnetization is $M_s = 900$ kA/m and the uniaxial anisotropy energy is $K_u = 657$ kJ/m$^3$ (see the Methods Section for a complete list of simulation parameters).
- **b**, Dot with saturation magnetization reduced to $M_s = 600$ kA/m at the bottom left edge.
- **c-e**, Dots with magnetic anisotropy energy reduced to $K_u = 550$ kJ/m$^3$ at the top left edge (**c**), entire circumference (**d**), and centre (**e**).

Homogeneous regions of the dots proceed almost unimpeded, with only minimal delay relative to the ideal dot. Other defect configurations, such as a continuous edge or a central spot with lower magnetic anisotropy, present more elaborate switching patterns (**Fig. 4.13d and e**), similar to those observed in **Fig. 4.8**. For all these defects, SOT-induced switching is fast, even though nucleation may start from a defect rather than the ideal edge side.
5 Asymmetric current-induced domain wall motion

Early observations of CIDWM in ultrathin HM/FM wires with out-of-plane magnetization evidenced two prominent effects: a large current-induced DW velocity $v_{DW}$ and a tilting of the DW relative to the propagation direction. The large $v_{DW}$ was initially attributed to a large FL torque stabilizing the DW magnetization against the Walker breakdown assuming Bloch DW configuration. Later it was recognized that due to the DMI the DWs are of Néel type (see Section 2.1.6) and that the large $v_{DW}$ is caused by the action of the DL SOT on the in-plane magnetization component of the DW. The tilt of the DW was initially observed in Pt/Co/Ni/Co wires by imaging the magnetic domains after a sequence of current pulses using MOKE microscopy and was later reproduced by micromagnetic simulations. Figure 5.1a illustrates the symmetry of the observed DW tilt for the four combinations of current and DW configuration, and the propagation direction of the DW front (green arrows). The micromagnetic simulations revealed that the DW tilt is a manifestation of the DMI in such wires. During the current-induced propagation the DW magnetization is deviated by the SOTs from its equilibrium configuration towards $\pm y$. However, the DMI tries to maintain the Néel DW configuration and thereby induces a rotation of the DW front. Recent theoretical work also shows that the DMI strength can be determined by measuring the static DW tilt as a function of a transverse field $B_y$.

The comparison with the DW tilt observed by the time- and spatially resolved switching experiments (Fig. 5.1b) discussed in the last chapter, reveals that the DW tilt differs by $\approx 90^\circ$ for a given current polarity and DW configuration. We argued that the apparent discrepancy comes from the time-resolved nature of the experiment in which a dynamic DW tilting process is measured, which can be different from the DW tilt observed at steady state. Furthermore, we argued that it can also be a manifestation of the $B^{FL} \propto \pm y$, which can alter the DW magnetization configuration and thereby influences the tilt angle and $v_{DW}$. Additionally, DW displacement experiments evidenced that $v_{DW}$ can be asymmetric in curved wires and y-shaped junctions, which is believed to be related to the DW tilting. Hence, a more profound understanding of the microscopic effects leading to the DW tilting and the observed asymmetry of
$v_{DW}$ will be beneficial for an improved SOT-induced switching efficiency and will allow for an increased DW motion in racetrack structures.

The aim of this chapter is to investigate the connection of $v_{DW}$ and the DW tilt in more detail in order to understand the microscopic mechanism that favours DW motion. We start by presenting a micromagnetic model of the DW tilt in 100 nm wide out-of-plane magnetized stripes.

**Domain wall tilting** In Fig. 5.2 we show micromagnetic simulations (see Section 3.5) of a 100 nm wide Pt/Co/AlO$_x$ stripe with the left (right) side being in the up (down) state. Figure 5.2a illustrates the equilibrium DW configuration, which is a left-handed Néel wall. The application of a transverse field $B_y$ induces the DW tilt (Fig. 5.2b). $B_y$ rotates the DW moments away from the effective DMI field towards a Bloch configuration, which causes a tilting of the DW in order to maintain the energetically favoured Néel configuration. The tilt is limited by the DW energy $E_{DW}$ due to an elongation of the DW with increasing tilt angle $\theta$. If instead of $B_y$, a positive current is injected into the stripe (we consider only the DL torque for the moment), the resulting DW tilt reverses and the DW starts propagating (Fig. 5.2c). In order to understand this dynamic DW tilt we consider $\frac{dm}{dt} \propto -T^{DL} - m \times T^{DL}$ (compare Eq. (2.16)). Hence, the DW magnetization is deviated towards $-y$ and $+z$ by the DL torque, as we show in the scheme of Fig. 5.2c. This dynamic process leads to the observed

![Figure 5.1 | Domain wall tilting.](image) Current-induced DW tilting symmetry for the four combinations of current and DW configuration measured by a, steady state by MOKE microscopy and b, by time-resolved x-ray measurement (see Chapter 4). In both cases we indicate the current polarity by the orange arrows and the propagating DW front by the green arrows.
propagation and tilting of the DW. Hence, the DW tilt angle depends on the competition between $T_{FL}$, $T_{DL}$, DMI and $E_{DW}$ as we show in Fig. 5.2d. For a constant $T_{DL}$ we can tune the DW tilt by changing the amplitude of $T_{FL}$, or equivalently by applying a bias field $B_y$.

**Asymmetric domain wall velocity** In order to investigate the relationship between the DW tilt angle and $v_{DW}$, we simulate the dynamics of an asymmetric

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**Figure 5.2 | Micromagnetic simulations of the domain wall tilt.** a, Equilibrium DW configuration of a 100 nm wide stripe. b, A magnetic field $B_y = 20$ mT rotates the DW magnetization towards $+y$ and induces thereby a static DW tilt. c, The injection of $j = 2 \cdot 10^8$ A cm$^{-2}$ with $T_{DL} = 18$ mT and $T_{FL} = 0$ mT per $10^8$ A cm$^{-2}$ induces a dynamic DW tilt and propagation. The lower schemes illustrate $dm/dt$ (according to Eq. (2.16)) due to $T_{DL}$ and the resulting DW tilt. d, The dynamic DW tilt can be tuned by changing the relative amplitude of $T_{FL}$ compared to $T_{DL} = 18$ mT per $10^8$ A cm$^{-2}$.
DW consisting of three straight sections. As we show in Fig. 5.3a, initially the magnetization in the left (right) side of the structure is pointing along $+z$ ($-z$), which leads to three Néel walls with tilt $\psi = -45^\circ$, $0^\circ$ and $45^\circ$ (see legend of Fig. 5.3c). For simplicity we neglect the FL torque, do not apply any magnetic field and inject only positive current pulses along the x-direction (electrons flowing to the left). All the other parameters are as introduced in Section 3.5.

Successive snapshots of the magnetic configuration plotted at intervals of 0.2 ns from each other (Fig. 5.3a) reveal that the different DW components prop-

Figure 5.3 | Asymmetric current-induced domain wall motion. a, Initial magnetization configuration of a 1.5 µm large square and snapshots of the magnetic configuration during the injection of a positive current of amplitude $j = 4.5 \cdot 10^8$ A cm$^{-2}$. b, Differential images obtained for various current amplitudes by subtracting two snapshots from each other with $\Delta t = 0.4$ ns. c, Evaluation of the DW velocity $v_n$, normal to the DW, as a function of current amplitude. The inset shows the definition of the different $v_{DW}$ components. d, Asymmetry of the CIDWM, defined as the ratio $v_n(\pm 45^\circ)/v_n(45^\circ)$, as a function of current amplitude.
Figure 5.4 | Domain wall components. a, Initial magnetization configuration illustrated by the individual components \( m_x, m_y \) and \( m_z \) (from left to right). The scheme to the right illustrates the DW magnetization. Due to the DMI, left-handed Néel walls are stabilized. 

b, Domain wall configuration during the injection of a current pulse with an amplitude \( j = 1.0 \times 10^8 \text{ A cm}^{-2} \). The DW magnetization is slightly deviated towards \(-y\). c, Domain wall configuration during the injection of a current pulse with an amplitude \( j = 4.5 \times 10^8 \text{ A cm}^{-2} \). The DW magnetization is strongly deviated towards \(-y\).

agate with distinct velocities for a current pulse of amplitude \( j = 4.5 \times 10^8 \text{ A cm}^{-2} \). If we imagine an elongated wire, the final domain state would be one with the DW normal being tilted at +45° with respect to the current-direction and the faster components (0°, −45°) will have disappeared.

Figure 5.3b shows differential images obtained by subtracting two snapshots separated by \( \Delta t = 0.4 \text{ ns} \) from each other, for multiple current densities. For each DW component we define the propagation velocity perpendicular to the DW \( v_n(45^\circ), v_n(0^\circ) \) and \( v_n(-45^\circ) \), as indicated in the legend of Fig. 5.3c. The values of \( v_n \) as a function of injected current density are shown in Fig. 5.3c. We find that \( v_n \) increases almost linearly with \( j \) for all components, however, with very distinct slopes. The asymmetry of the CIDWM, which we define as the ratio \( v_n(-45^\circ)/v_n(45^\circ) \) increases strongly as we increase the injected current pulse amplitude (Fig. 5.3h).
Micromagnetic model  In the following we elucidate the process that leads to the asymmetric CIDWM. Figure 5.4a shows the three magnetization components \( m_x, m_y \) and \( m_z \) (from left to right) of the initial equilibrium configuration. The scheme to the right is an illustration of the DW magnetization configuration. In agreement with the DMI\(^{50,52,149}\) we observe that left-handed Néel-walls are stabilized. In Fig. 5.4b we illustrate the effect of a current pulse with moderate amplitude (\( j = 1 \cdot 10^8 \text{ A cm}^{-2} \)) on the DW magnetization. With the same argumentation used before we understand that \( T_{DL} \) rotates the DW magnetization, against the DMI, slightly towards \(-y\) and hence, \( m_x \) is largest (smallest) for the \( 0^\circ \) (45\(^\circ\)) DW. Further, \( T_{DL} \propto m_x \) due to the DL torque symmetry (\( T_{DL} = T_{DL} m_x \times (y \times m) \)) consistent with the \( v_{DW} \) being the largest (smallest) for the \( 0^\circ \) (45\(^\circ\)) DW. The injection of large current densities deviates the DW magnetization substantially from its equilibrium position towards \(-y\), which results in very different \( m_x \) for the three DW components and hence, in a pronounced asymmetry of the CIDWM, as illustrated in Fig. 5.4c. Alternatively, the increased \( v_{DW} \) of the \(-45^\circ\) (and the \( 0^\circ\)) DW can be understood in terms of energy. The energy of the system is minimized if the DWs are of Néel type and hence, the DWs which are closer to a Bloch wall type propagate faster in order to reduce the DW energy by increasing the length of the energetically favoured Néel walls. We remark that the final dynamic tilt angle can be different from the initial 45\(^\circ\). However, the emergence of the final tilt happens on a longer time
scale, which increases with the stripe width as pointed out by Boulle et al.\textsuperscript{166}.

**Field-like torque**  Next, we introduce the FL torque, or any equivalent in-plane field pointing along $+y$ and study the asymmetric CIDWM. $B^{FL}$ counteracts the DL torque induced DW magnetization rotation towards $-y$ and thereby increases the $m_x$ components of all three DWs. Consistent with our previous argumentation we find that the DW velocity of all three DW components increases and the asymmetry decreases. This is consistent with our experimental findings that reversal is assisted by applying an in-plane field $B_y$ in the direction of $B^{FL}$ (see Section 4.1.2) and the report of Emori et al.\textsuperscript{50} in which they showed that the current-induced DW velocity increases for a $B_y$ applied in the direction of $B^{FL}$ (see Fig. 2.9c).

**Conclusions**  Similar to crystal growth, where the facet with the slowest growth-rate determines the final crystal shape, during CIDWM we expect the final DW tilt to be governed by the slowest DW component. This implies that CIDWM studies in which the DW tilt is probed after the application of a sequence of current pulses, measure the rest state of the domain walls\textsuperscript{157,162,163,166} which can be different from the dynamic DW tilt governed by the fastest DW component\textsuperscript{79,183}. We point out that the fastest DW component reported by Safeer et al.\textsuperscript{183} in Pt/Co bilayers is consistent with our micromagnetic simulations, as well as the DW propagation direction we measured by STXM (compare the lower right case in Fig. 5.1b with Fig. 5.4c). Additionally, we emphasize that the DW tilt observed at steady state\textsuperscript{162} does indeed correspond to the slowest DW component reported above.

The findings presented in this chapter pave the way for a better understanding of SOT-induced DW dynamics and related effects such as DW tilting. The improved understanding of the microscopic interplay of the DMI, SOTs and external fields on DW motion will allow to increase the DW speed in racetrack devices and improve the SOT switching efficiency of single and multi-bit MRAM devices.
6 Conclusions and Outlook

In the last years the understanding of nonequilibrium spin-orbit coupling effects in ultrathin HM/FM bilayers has improved at a very high pace. Among such effects, current-induced spin-orbit torques have shown outstanding possibilities to drive fast magnetization switching and to control DW motion in HM/FM heterostructures.

In this thesis we focused on the SOT-induced magnetization reversal of perpendicularly magnetized Pt/Co dots and on the asymmetry of current-induced domain wall motion in out-of-plane magnetized HM/FM stripes. The combination of time- and spatially resolved measurements with micromagnetic simulations led to a deeper understanding of the microscopic switching process and DW dynamics. Our findings are of interest for the development of future spinorbitronic applications as well as for fundamental understanding of the contributing effects. In the following, we summarize the main findings and give an outlook on future studies.

The time- and spatially-resolved SOT-induced magnetization switching measurements presented in this thesis, provide unprecedented insights into SOT-induced magnetization dynamics during the reversal process and the accompanying spatial evolution of the magnetization. Our study revealed that the switching of Pt/Co dots takes place by a current-induced DW nucleation process and subsequent DW propagation. We could demonstrate that the switching occurs during the current pulse and that the reversal speed increase as a function of current amplitude, and hence allows for sub-ns switching by adjusting the current pulse amplitude. Furthermore, as the switching unfolds along deterministic and reproducible paths, the extent of the reversal can be reliably controlled by fine tuning the duration and amplitude of the injected current pulse as well as by tailoring the sample dimension and geometry. Most interesting from an application point of view is the demonstration of reliable sub-ns magnetization switching for more than $10^{12}$ cycles without any meaningful incubation delay, which makes SOTs a serious competitor of STT-based spintronics application such as MRAMs. Additionally, we showed that the DW nucleation, as well as the DW propagation, has a four-fold symmetry which we attribute to the concurring action of the SOTs, the DMI and the external field. The DW propagates at an angle of $\approx 45^\circ$ with respect to the current direction, opposite to the DW tilt recently reported at steady state in racetrack structures. Additional all-
electrical pulsed measurements and micromagnetic simulations, uncovered the important role of the FL torque in assisting the DW dynamics and in determining the nucleation point in addition to the DMI and the DL torque. Thus, we anticipate that tuning the ratio $T_{FL}/T_{DL}$, or exploiting the Oersted field in well-engineered structures, may lead to more efficient SOT switching, beneficial for the development of SOT-MRAM.

The investigated asymmetry of the current-induced domain wall motion paves the way for DW propagation with increased efficiency by furthering the understanding of the dynamic processes defining the DW propagation speed. We presented a simple model describing the DW tilt in the rest state and during dynamic propagation based on the competing action of the DMI, magnetic bias-fields and SOTs. Furthermore, we showed that the same model can be applied in order to explain the asymmetric DW speed observed in wires with PMA. We find that the experimentally observed DW tilt in quasi static measurements corresponds to the slowest DW component and consistently, the fastest DW component corresponds to the propagating DW front during SOT-induced magnetization switching observed in the STXM measurements.

Several points remain outstanding and solicit for additional experiments and investigations. Below, we list some suggestions for new projects and experimental studies.

- We propose to investigate the effect played by the geometry of the magnetic dots. Micromagnetic simulations reveal that the reversal speed is strongly linked to the specific shape of the magnetic dots. The DW nucleation process is favoured in certain geometries and the internal DW structure is related to the width of the current line and may thereby change the switching mechanism. Figure 6.1 shows preliminary results of shape dependent magnetization switching in Co dots patterned into squares, diamonds (rotated square) and ellipses. Independent of the geometry, the switching happens through a DW nucleation and propagation process. However, in this batch of samples the nucleation seems to be dominated by magnetic defects and hence does not show the expected edge driven nucleation process.

- Complementary to the time- and spatially resolved switching experiments, the investigation of sub-threshold switching may enable the development of multi state memories. This could be facilitated by engineering magnetic structures with intentional pinning centres. One could imaging that a DW
Figure 6.1 | Shape-dependent magnetization switching. SOT-driven magnetization reversal in a, a 500 nm wide square and b, a diamond, c, a 500 nm × 200 nm ellipse and d, a 200 nm × 500 nm ellipse. In all cases the switching happens through a domain wall nucleation and propagation process. However, the nucleation seems to be dominated by magnetic defects and hence does not show the expected edge driven nucleation process.

Propagating through a magnetic structure would jump from pinning centre to pinning centre, thereby changing the readout voltage, either by means of TMR or Hall effect, in a discrete and predictable way depending on the length and amplitude of the injected current pulse.

- As described in the main text, one of the main handicaps of SOT-induced magnetization switching, compared to STT, is the need of an in-plane bias-field. Significant investigations have been carried out in order to eliminate this burden. We propose to extend the perpendicularly magnetized Co dot with an in-plane magnetized region which would act like a small permanent magnet and may result in a significant reduction of $B_x$. We could successfully demonstrate that it is possible to incorporate within the same magnetic microstructure in- and out-of-plane magnetized regions by selective oxidation of the Co/Al interface. Additionally, an inhomogeneous current distribution, originating from curved current lines, could be exploited in order to enable zero-field switching.
The development of all-electrical time-resolved switching measurements would be of fundamental as well as of practical interest. Such single shots measurements would enable extensive in-house studies as a function of field, current pulse amplitude and duration and will further strengthen the understanding of the reversal process. For example it would allow to study Ta/CoFeB based structures. Due to the element specificity of the STXM and the reduced Co (and Fe) volume in CoFeB structures, it is very challenging to investigate these bilayers by means of x-ray studies. The investigation of Ta/CoFeB bilayers is very interesting as negligible DMI and Bloch DWs were reported in this system. Furthermore, in contrast to the x-ray measurements conducted in the framework of this thesis, such single shot measurements would allow the investigation of stochastic switching effects. Additionally, the SOT-induced switching process in MTJs could directly be compared to the STT-induced switching process by time-resolved measurements of the TMR during reversal. Finally, the investigation could be extended to antiferromagnetic structures for which SOT-induced switching was recently demonstrated. The switching in such structures happens through the reversal of single grains and shows large aftereffects attributed to the relaxation of the grains, which makes it very interesting for time-resolved studies.
Appendices
A Calculus

A.1 Estimate of effective anisotropy field

The magnetization equilibrium position is given by the competition of the Zeeman energy and the anisotropy field and hence their corresponding magnetic torques have to cancel each other in equilibrium $T_Z + T_k = 0$. From Section 2.1.1 we know that a magnetic torque acting is defined as $T = m \times B$. Furthermore, we make use of the effective out-of-plane anisotropy field defined in Section 2.1.4 as $B_k = B_k (m \cdot z) \cdot z$. The equilibrium condition can thus be written as

$$T_k = m \times [B_k (m \cdot z) z] = -m \times B_{ext} = -T_Z. \quad (A.1)$$

Using spherical coordinates (Eq. (A.7)) for the magnetization, parametrized by the polar ($\theta$) and azimuthal ($\varphi$) angle, and the magnetic field, parametrized by the amplitude $B_{ext}$ and the polar ($\theta_B$) and azimuthal ($\varphi_B$) angle, we can solve above equation for $B_k$ and find

$$B_k = B_{ext} \left[ \frac{\sin(\theta_B)}{\sin(\theta)} - \frac{\cos(\theta_B)}{\cos(\theta)} \right]. \quad (A.2)$$

We utilized that $\varphi = \varphi_B$ in the case of uniaxial out-of-plane anisotropy.

A.2 Spin-orbit torques and effective fields

The definition of the SOTs for a positive current injected into Pt/Co/AlO$_x$ along the x-axis is:

$$T_{FL}^F = (m \times y) \left[ T_{0}^{FL} + T_{2}^{FL} (z \times m)^2 \right] - m \times (z \times m) (m \cdot x) T_{2}^{FL} \quad (A.3)$$

$$T_{DL}^F = T_{0}^{DL} [m \times (y \times m)] + T_{2}^{DL} (z \times m) (m \cdot x), \quad (A.4)$$

with $T_{i}^{DL(FL)} > 0$. Alternatively, we can express the SOTs using the definition of effective fields ($B = T \times m$). With $T_{2}^{FL} = T_{2}^{DL} = 0$ we obtain for $B^{FL}$ and $B^{DL}$

$$B^{FL} = B_{0}^{FL} y, \quad (A.5)$$
and

\[ \mathbf{B}_{DL}^{DL} = B_0^{DL} (\mathbf{y} \times \mathbf{m}) = \begin{pmatrix} m_z \\ 0 \\ -m_x \end{pmatrix} = \begin{pmatrix} \cos(\theta) \\ 0 \\ -\sin(\theta) \cos(\varphi) \end{pmatrix}, \] (A.6)

respectively. Using spherical coordinates

\[ \mathbf{m} = \begin{pmatrix} \sin(\theta) \cos(\varphi) \\ \sin(\theta) \sin(\varphi) \\ \cos(\theta) \end{pmatrix}, \] (A.7)

and the corresponding spherical unit vectors

\[ \mathbf{e}_r = \begin{pmatrix} \sin(\theta) \cos(\varphi) \\ \sin(\theta) \sin(\varphi) \\ \cos(\theta) \end{pmatrix}, \quad \mathbf{e}_\theta = \begin{pmatrix} \cos(\theta) \cos(\varphi) \\ \cos(\theta) \sin(\varphi) \\ -\sin(\theta) \end{pmatrix}, \quad \mathbf{e}_\varphi = \begin{pmatrix} -\sin(\varphi) \\ \cos(\varphi) \\ 0 \end{pmatrix}. \] (A.8)

The polar and azimuthal components of the individual SOTs are evaluated to

\[ B_{\theta}^{FL} = \mathbf{B}_{\theta}^{FL} \cdot \mathbf{e}_\theta = B_0^{FL} \cos(\theta) \sin(\varphi), \] (A.9)

\[ B_{\varphi}^{FL} = \mathbf{B}_{\varphi}^{FL} \cdot \mathbf{e}_\varphi = B_0^{FL} \cos(\varphi), \] (A.10)

\[ B_{\theta}^{DL} = \mathbf{B}_{\theta}^{DL} \cdot \mathbf{e}_\theta = B_0^{DL} \cos(\varphi), \] (A.11)

\[ B_{\varphi}^{DL} = \mathbf{B}_{\varphi}^{DL} \cdot \mathbf{e}_\varphi = -B_0^{DL} \sin(\varphi) \cos(\theta). \] (A.12)

Hence, the polar and azimuthal components of the current induced effective fields are given by

\[ B_\theta^{I} = B_\theta^{FL} + B_\theta^{DL} = B_0^{FL} \cos(\theta) \sin(\varphi) + B_0^{DL} \cos(\varphi), \] (A.13)

and

\[ B_\varphi^{I} = B_\varphi^{FL} + B_\varphi^{DL} = B_0^{FL} \cos(\varphi) - B_0^{DL} \sin(\varphi) \cos(\theta). \] (A.14)

With the same argumentation we can obtain the polar and azimuthal components of the current induced effective in the general case \( T_2^{FL} \neq T_2^{DL} \neq 0: \)

\[ B_\theta^{I} = \cos(\theta) \sin(\varphi) \left[ B_0^{FL} + B_2^{FL} \sin^2(\theta) \right] \]
\[ + \left[ B_0^{DL} + B_2^{DL} \sin^2(\theta) \right] \cos(\varphi), \] (A.15)

and

\[ B_\varphi^{I} = B_0^{FL} \cos(\varphi) - B_0^{DL} \cos(\theta) \sin(\varphi). \] (A.16)
A.3 Approximation of harmonic signals

In the following section we show the computation of the first and second harmonic Hall signals due to small ac current $I(t) = I_0 \cos(\omega t)$ injection (Section 3.3). For brevity we limit the discussion to the transverse Hall signal $V_H = R_H(t) I_0 \cos(\omega t)$. $R_H$ depends on the magnetization position through $R_{AHE}$ and $R_{PHE}$. Hence we can write

$$R_H(t) = R_H \left[ B_0 + B^I(t) \right].$$  \hfill (A.17)

Here, $B^I(t)$ parametrizes the current induced fields (SOTs and Oersted field) and $B_0$ the static magnetic fields ($B_{\text{ext}}$ and $B_k$). As $B^I(t) \propto I(t)$ we can approximate $R_H(t)$ as

$$R_H(t) \approx R_H(B_0) + \frac{dR_H}{dB^I} \cdot B^I \cos(\omega t).$$  \hfill (A.18)

Using in terms of total transverse voltage we get

$$V_H(t) = I_0 \cos(\omega t) \left[ R_H(B_0) + \frac{dR_H}{dB^I} \cdot B^I \cos(\omega t) \right]$$

$$= I_0 \left[ R^0_H + R^\omega_H \cos(\omega t) + R^{2\omega}_H \cos(2\omega t) \right],$$  \hfill (A.19)

with the harmonics

$$R^0_H = R^2_H = \frac{1}{2} \frac{dR_H}{dB^I} \cdot B^I, \quad R^\omega_H = R_H(B_0).$$  \hfill (A.20)

We used the identity $\cos^2(\omega t) = \frac{1}{2} [1 + \cos(2\omega t)]$. We compute $R^{2\omega}_H$ utilizing the definition of the first harmonic signal introduced in Section 3.1,

$$R^{2\omega}_H = R_{AHE} \cos(\theta) + R_{PHE} \sin(2\varphi) \sin^2(\theta),$$  \hfill (A.21)

and by calculating the derivative, which results in

$$R^{2\omega}_H = \frac{1}{2} \frac{dR_H}{dB^I} \cdot B^I$$

$$= \frac{1}{2} \left[ R_{AHE} - 2R_{PHE} \cos(\theta) \sin(2\varphi) \right] \frac{d\cos(\theta)}{dB^I} \cdot B^I$$

$$+ \frac{1}{2} 2R_{PHE} \sin^2(\theta) \frac{d\sin(2\varphi)}{dB^I} \cdot B^I.$$

(A.22)
A.4 Spin-orbit torque analysis

As we are only sensitive to the polar and azimuthal component of \( B_I \) we can simplify the two derivatives appearing in the previous equation as

\[
\frac{d \cos(\theta)}{dB^I_\theta} B^I_\theta = \frac{d \cos(\theta)}{dB^I_\theta} B^I_\theta,
\]

(A.23)

and

\[
\frac{d \sin(2\varphi)}{dB^I_\varphi} B^I_\varphi = \frac{d \sin(2\varphi)}{dB^I_\varphi} B^I_\varphi.
\]

(A.24)

We can further simplify the expression by using the following substitutions applicable for field scans (variable \( B_{\text{ext}} \), constant \( \theta_B \) and \( \varphi_B \))

\[
\frac{d \cos(\theta)}{dB^I_\theta} B^I_\theta = \frac{d \cos(\theta)}{dB^I_\theta} \frac{1}{\sin(\theta_B - \theta)} B^I_\theta
\]

(A.25)

\[
\frac{d \sin(2\varphi)}{dB^I_\varphi} B^I_\varphi = \frac{d \sin(2\varphi)}{dB^I_\varphi} 2 \cos(2\varphi) \frac{d \varphi}{dB^I_\varphi} \approx \frac{2 \cos(2\varphi)}{B_{\text{ext}} \sin(\theta_B)} B^I_\varphi.
\]

(A.26)

Finally, we obtain the second harmonic signal

\[
R^{2\omega}_{H} = \frac{1}{2} [R_{AH}^E - 2R_{PH}^E \cos(\theta) \sin(2\varphi)] \frac{d \cos(\theta)}{dB^I_\theta} \frac{1}{\sin(\theta_B - \theta)} B^I_\theta
\]

\[
+ \frac{1}{2} R_{PH}^E \sin^2(\theta) \frac{2 \cos(2\varphi)}{B_{\text{ext}} \sin(\theta_B)} B^I_\varphi
\]

(A.27)

where \( B^I_\theta \) and \( B^I_\varphi \) are defined in Eq. [A.15] and [A.16], respectively.

A.4 Spin-orbit torque analysis

In Section 3.3.2 of the main text we introduced the matrix method for the SOT analysis based on two independent measurements. In the following section we expand the reasoning to four independent measurements, which gives access to \( B^I_\theta \), as well as to \( B^I_\varphi \).

Starting from Eq. [A.27] and using the definition of \( B^I_\theta \) and \( B^I_\varphi \) given in Eq. [A.15] and [A.16], we can write the matrix equation which has to be solved

\[
\begin{pmatrix}
R^{2\omega}_{H,0} \\
R^{2\omega}_{H,45} \\
R^{2\omega}_{H,90} \\
R^{2\omega}_{H,315}
\end{pmatrix} =
\begin{pmatrix}
A_0 & B_0 & C_0 & D_0 \\
A_{45} & B_{45} & C_{45} & D_{45} \\
A_{90} & B_{90} & C_{90} & D_{90} \\
A_{315} & B_{315} & C_{315} & D_{315}
\end{pmatrix}
\begin{pmatrix}
B^D_L \\
B^D_R \\
B^F_L \\
B^F_R
\end{pmatrix} +
\begin{pmatrix}
R^{2\omega}_{H,0} \\
R^{2\omega}_{H,45} \\
R^{2\omega}_{H,90} \\
R^{2\omega}_{H,315}
\end{pmatrix}
\]

(A.28)
A.4. Spin-orbit torque analysis

From a simple comparison we can define the coefficients

\[ A_i = \zeta_i \cos(\varphi_i), \]  \hspace{1cm} (A.29)

\[ B_i = \zeta_i \cos(\theta) \sin(\varphi_i), \]  \hspace{1cm} (A.30)

\[ C_i = \eta_i \cos(\theta) \sin(\varphi_i), \]  \hspace{1cm} (A.31)

\[ D_i = \eta_i \cos(\varphi_i), \]  \hspace{1cm} (A.32)

here, \( i = 0^\circ, 45^\circ, 90^\circ \) and \( 315^\circ \). We defined \( \zeta_i \) and \( \eta_i \) as

\[ \zeta_i = \frac{1}{2} \left[ R_{AHE} - 2 R_{PHE} \cos(\theta) \sin(2\varphi) \right] \frac{d \cos(\theta)}{d B_{ext}} \frac{1}{\sin(\theta_B - \theta)}, \]  \hspace{1cm} (A.33)

and

\[ \eta_i = \frac{1}{2} R_{PHE} \sin^2(\theta) \frac{2 \cos(2\varphi)}{B_{ext} \sin(\theta_B)}, \]  \hspace{1cm} (A.34)

respectively, which is the first and second term of Eq. (A.27). \( R_{\omega \theta}^{2\omega \theta} \) is the thermal contribution to \( R_{\omega \theta}^{2\omega \theta} \) which has to be considered and evaluated separately (compare Eq. (3.11)). At each measurement point \( R_{\omega \theta}^{2\omega \theta} \), Eq. (A.28) is evaluated and solved for the SOTs which are parametrized by \( B_\theta^I \) and \( B_\varphi^I \).
B  Comparison between filtered and raw data

Given the very low photon count rate at the detector position and the relatively low x-ray absorption contrast of the 1 nm thick Co layer, the acquisition time of a single sequence of images with 100 ps resolution requires several hours of integration. In order to enhance the magnetic contrast, the intensity of each frame in Figs. 4.4, 4.7 and 4.8 has been normalized to the non-magnetic background. Moreover, to improve the signal-to-noise ratio, the intensity of each pixel has been low-pass filtered ($f_c \approx 10$ MHz) along the time axis. We have checked that the filtering process does not introduce a significant distortion of the data. For completeness, we compare the raw and filtered images side-by-side in Figs. B.1, B.2 and B.3.

![Comparison of the raw and filtered data shown in Fig. 4.4.](image)

Figure B.1  |  Comparison of the raw and filtered data shown in Fig. 4.4.  

a, Raw and b, low-pass filtered images.
Figure B.2 | Comparison of the raw and filtered data shown in Fig. 4.7.  
a, Raw and b, low-pass filtered images.

Figure B.3 | Comparison of the raw and filtered data shown in Fig. 4.8.  
a, Raw and b, low-pass filtered images.
C Analysis programs

C.1 Fitting procedure for spin-orbit torques measurement

The MATLAB script below analyses harmonic Hall voltage measurements of out-of-plane magnetized samples using the fitting procedure described in Section 3.3.2 in which the second harmonic signal is fitted by Eq. (3.9). The script requires input parameters $\theta_B$ and $R_{PHE}$ and can simultaneously fit measurements with different $\varphi_B$.

```
%%%% Parameter input %%%%%
R_PHE = 0.02;
theta_B_set = 86;

%%%% read and prepare data %%%%%
[fileName, readPath, numberOfFiles] = getFilesToRead();
output = [];
for fileNumber = 1 : numberOfFiles
    name = char(fileName(fileNumber));
data = importFile(readPath, name);
% getting phi_B from file name %
phi_B_set = getPhiFromFileName(name);
disp(['phi_B = ', num2str(phi_B_set)]);
% interpolating data linearly%
[dataInterpl] = interpolateData(data);
% centering data %
out = centeringData(dataInterpl);
field = out(:, 1);
R1f = out(:, 2);
R2f = out(:, 3);
% ploting raw data %
figure()
plot(field, R1f, 'o')
figure()
plot(field, R2f, 'o')
% get R_AHE %
[PHEcontrib, AHEcontrib] = symmetrizeData(field, R1f);
R_AHE = max(AHEcontrib);
AHEcontrib(abs(AHEcontrib(:)) > R_AHE ) = R_AHE;
% calculate theta(B_ext) %
theta = acos(AHEcontrib / R_AHE);
data = [field, R1f, R2f, theta];
% remove thermal contribution %
[data] = removeANE(data);
```
C.1. Fitting procedure for spin-orbit torques measurement

% define region of interest to be analyzed %
[data1, data2] = cutData(data, 'maxToZeroField');

% setting up fit procedure %
angles = calcAngles(data1, data1(:, 4), phi_B_set, theta_B_set, 'deg');
dcos_dfield = dfx_dx(data1(:, 1), cos(data1(:, 4)));
dataToAdd = {data1(:, 1), data1(:, 3), dcos_dfield, angles, R_AHE, R_PHE};

output = arrangeDataForSOTfit(dataToAdd, output);

```
estimate = nlinfit(output, output(:, 2), @multiR2fFit, [10, 10, 10, 1e-3])
[R2f] = multiR2fFit(estimate, output);
```

%%%%%%%%% plot output %%%%%%%%%
id = output(:, 9);
figure()
gscatter(output(:, 1), output(:, 2), id)
hold all
for IDs = 1:output(end, 9)
    plot(output((id == IDs), 1), R2f(id == IDs), 'lineWidth', 3);
end

% rearrangeDataForSOTfit

function [output] = rearrangeDataForSOTfit(dataToAdd, existingData)
% dataToAdd will be appended to existingData, if existingData=[]
% data will be created accordingly
% dataToAdd should be a cellArray with field, R2f, dcos_dfield, theta, phi, theta_B, R_AHE, R_PHE

lengthData = length(cell2mat(dataToAdd(1)));
scratch = ones(lengthData, 1);

```
tempOutput(:, 1) = cell2mat(dataToAdd(1)); %field
tempOutput(:, 2) = cell2mat(dataToAdd(2)); %R2f
tempOutput(:, 3) = cell2mat(dataToAdd(3)); %dcos_dfield
angles = cell2mat(dataToAdd(4));
tempOutput(:, 4) = angles(:, 1); %theta
tempOutput(:, 5) = angles(:, 2); %phi
tempOutput(:, 6) = angles(:, 3); %theta_B
tempOutput(:, 7) = scratch * cell2mat(dataToAdd(5)); %R_AHE
tempOutput(:, 8) = scratch * cell2mat(dataToAdd(6)); %R_PHE
```

if length(existingData) < 9 % should test that existingData is empty -> alternative with input parsing
    output = tempOutput;
    output(:, 9) = 1;
else
    tempOutput(:, 9) = scratch * (existingData(end, 9) + 1); % updating fitID
    output = [existingData; tempOutput];
end
C.2 Spin-orbit torques from matrix

Below we show a MATLAB script which analyses field scan measurements based on the matrix method described in the main text (Section 3.3.2) and in Appendix A.4 for four independent measurements. The script requires the input parameters $\theta_B$, $R_{PHE}$ and $R_{2\Omega}^{\omega}\nabla T = cTh_z$ which has to be obtained from an additional measurement at $\varphi_B = 45^\circ$ and $\theta_B = 90^\circ$ in the high field limit (see Eq. (3.11)).

```matlab
function [ R2f ] = multiR2fFit( estimate, dataToFit )
% Definition of the second harmonic to be fitted
%
% taken into account are: B_FL0 B_FL2 B_DL c_Th
% dataToFit = [field, R2f, dcos_dfield, theta, phi, theta_B, R_AHE, R_PHE];

field = dataToFit (: , 1);
dcos_dfield = dataToFit (: , 3);
theta = dataToFit (: , 4);
phi = dataToFit (: , 5);
theta_B = dataToFit (: , 6);
R_AHE = dataToFit (: , 7);
R_PHE = dataToFit (: , 8);
B_FL0 = estimate(1);
B_FL2 = estimate(2);
B_DL = estimate(3);
c_Th = estimate(4);

A = ( R_AHE - 2 * R_PHE .* cos(theta).* sin(2* phi)).* dcos_dfield ./ sin(theta_B -theta);
B = 2 * R_PHE .* sin(theta).^2 .*cos(2 * phi) ./ ( field .* sin(theta_B));
C = cos(theta).*sin(phi);
D = cos(phi);
E = sin(theta) .* sin(theta);
F = sin(theta) .* cos(phi);

% comment this line if you want to get c_Th from the fit, if you
% know c_Th uncomment and set the value here
%c_Th = 0*7.57561e-5;
R2f = -sign(field) .* (A.*(B_FL0 + B_FL2.*E).*C + B_DL.*D) + B .* (B_FL0.*D - B_DL.*C)) + c_Th.*F;
end

C.2 Spin-orbit torques from matrix

Below we show a MATLAB script which analyses field scan measurements based on the matrix method described in the main text (Section 3.3.2) and in Appendix A.4 for four independent measurements. The script requires the input parameters $\theta_B$, $R_{PHE}$ and $R_{2\Omega}^{\omega}\nabla T = cTh_z$ which has to be obtained from an additional measurement at $\varphi_B = 45^\circ$ and $\theta_B = 90^\circ$ in the high field limit (see Eq. (3.11)).

% Parameter input %
theta_B = 86;
R_PHE = 0.02;
```
\[ c_{Th_z} = 6.897 \times 10^{-5}; \quad \% \text{Thermal contribution} \]

\begin{verbatim}
phi = 0;
theta_B_const = theta_B * pi / 180;
selectAngle = num2str(phi);
[file, readPath] = uigetfile([defaultPath,'/*.*'], ['Select file phi = ', selectAngle, ' ']);
values0 = importFile(readPath, fileName);
% interpolating data, removing ANE contribution and centering data
[cleanedData, R_AHE_0, R_ANE_0, iterField] = cleaningData(values0);
Bext_0 = cleanedData(:, 1);
R1f_0 = cleanedData(:, 2);
R2f_0 = cleanedData(:, 3);
phi_const = phi * pi / 180;
theta_B = pi / 2 * (1 - sign(Bext_0)) + theta_B_const * sign(Bext_0);
[DL_theta_0, FL_theta_0, DL_phi_0, FL_phi_0, phi_0, theta_0, Bk_0, sym1f_0] = determineCoeff(Bext_0, theta_B, phi_const, R_PHE, R1f_0);

phi = 45;
selectAngle = num2str(phi);
[file, readPath] = uigetfile([readPath,'/*.*'], ['Select file phi = ', selectAngle, ' ']);
importFileName = [readPath, fileName];
values45 = importFile(readPath, fileName);
[cleanedData, R_AHE_45, R_ANE_45] = cleaningData(values45, Bext_max, dB_ext);
Bext_45 = cleanedData(:, 1);
R1f_45 = cleanedData(:, 2);
R2f_45 = cleanedData(:, 3);
phi_const = phi * pi / 180;
theta_B = pi / 2 * (1 - sign(Bext_45)) + theta_B_const * sign(Bext_45);
[DL_theta_45, FL_theta_45, DL_phi_45, FL_phi_45, phi_45, theta_45, Bk_45, sym1f_45] = determineCoeff(Bext_45, theta_B, phi_const, R_PHE, R1f_45);

phi = 90;
selectAngle = num2str(phi);
[file, readPath] = uigetfile([readPath,'/*.*'], ['Select file phi = ', selectAngle, ' ']);
values90 = importFile(readPath, fileName);
[cleanedData, R_AHE_90, R_ANE_90] = cleaningData(values90, Bext_max, dB_ext);
Bext_90 = cleanedData(:, 1);
R1f_90 = cleanedData(:, 2);
R2f_90 = cleanedData(:, 3);
\end{verbatim}
phi_const = phi * pi / 180;
theta_B = pi / 2 * (1 - sign(Bext_90)) + theta_B_const * sign(Bext_90);
[DL_theta_90, FL_theta_90, DL_phi_90, FL_phi_90, phi_90, theta_90, Bk_90, sym1f_90] = determineCoeff(Bext_90, theta_B, phi_const, R_PHE, Rif_90);

%%%%%%%%%%%%%%%% phi = 315 %%%%%%%%%
phi = 315;
Bext_315 = -Bext_45;
R1f_315 = -R1f_45;
R2f_315 = -R2f_45;
scanLength = length(Bext_45);
scanPointNumber = (scanLength+1)/2;
Bext_315 = [Bext_315(scanPointNumber : scanLength); Bext_315(2 : scanPointNumber)];
R1f_315 = [R1f_315(scanPointNumber : scanLength); R1f_315(2 : scanPointNumber)];
R2f_315 = [R2f_315(scanPointNumber : scanLength); R2f_315(2 : scanPointNumber)];
phi_const = phi * pi / 180;
theta_B = pi / 2 * (1 - sign(Bext_315)) + theta_B_const * sign(Bext_315);
[DL_theta_315, FL_theta_315, DL_phi_315, FL_phi_315, phi_315, theta_315, Bk_315, sym1f_315] = determineCoeff(Bext_315, theta_B, phi_const, R_PHE, Rif_315);

%%%%%%%%%%%%%%%% solving matrix equation %%%%%%%%%
R2f_output = zeros(length(Bext_0), 4);
for i = 1:length(Bext_0)
matrix = [DL_theta_0(i), FL_theta_0(i), DL_phi_0(i), FL_phi_0(i);
DL_theta_45(i), FL_theta_45(i), AdPhi45(i), FL_phi_45(i);
DL_theta_90(i), FL_theta_90(i), DL_phi_90(i), FL_phi_90(i);
DL_theta_315(i), FL_theta_315(i), DL_phi_315(i), FL_phi_315(i)];
R2f = [R2f_0(i) - cTh_z*sin(theta_0(i))*cos(phi_0(i));
R2f_45(i) - cTh_z*sin(theta_45(i))*cos(phi_45(i));
R2f_90(i) - cTh_z*sin(theta_90(i))*cos(phi_90(i));
R2f_315(i) - cTh_z*sin(theta_315(i))*cos(phi_315(i))];
SOTs(i,:) = linsolve(matrix, R2f);
R2f_output(i,:) = R2f';
end

%%%%%%%%%%%%%%%% create output array %%%%%%%%%
DL_theta = SOTs(:,1);
FL_theta = SOTs(:,2);
DL_phi = SOTs(:,3);
FL_phi = SOTs(:,4);
sin20 = sin(theta_0).*sin(theta_0);
sin245 = sin(theta_45).*sin(theta_45);
sin290 = sin(theta_90).*sin(theta_90);
sin2 = (sin20 + sin245 + sin290)/3;
C.2. Spin-orbit torques from matrix 100

\[
B_k = \frac{1}{3} \left( B_{k_0} + B_{k_45} + B_{k_90} \right);
\]

\[
\text{output} = [\sin^2, \text{abs}(DL_{\theta}), \text{abs}(FL_{\theta}), \text{abs}(DL_{\phi}), FL_{\phi}, B_k, \text{sym1f}_{45}];
\]

\[
R_{\text{AHE}} = \frac{1}{3} \left( R_{\text{AHE}_0} + R_{\text{AHE}_{45}} + R_{\text{AHE}_{90}} \right);
\]

\[
\text{disp('AHE:')};
\]

\[
R_{\text{ANE}} = \frac{1}{3} \left( R_{\text{ANE}_0} + R_{\text{ANE}_{45}} + R_{\text{ANE}_{90}} \right);
\]

\[
\text{disp('ANE:')};
\]

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D Macrospin simulations

D.1 Harmonic Hall voltage simulation

The following MATLAB code simulates harmonic Hall voltage measurements in samples with PMA. The code is based on the harmonic ac-measurements introduced in Section 3.3 It outputs $R_H^e$ and $R_H^{2\omega}$ as a function of applied external field according to Eqs. (3.1) and (3.9).

\[
\begin{align*}
\text{%%%%%%%%% Parameter input %%%%%%%%%}
\end{align*}
\]

\[
\begin{align*}
\text{% external field definition} \\
\theta_B &= 86; \quad \text{in degress} \\
\phi_B &= 45; \quad \text{in degess} \\
\text{% in mT, format [B1, B2, dB]} \\
\text{fieldSetArray} &= [2000, 1200, 10; 1200, -1200, 5; -1200, -2000, 10]; \\
\text{hysteresis} &= \text{true}; \quad \text{% makes a hysteresis loop} \\
\text{addNoise} &= \text{false}; \quad \text{% adds noise to the field} \\
\text{% Magnetic properties in mT} \\
B_k1 &= 800; \\
B_k2 &= 0; \\
B_c &= 500; \\
\text{% SOT’s in mT per 10^8 A cm^-2} \\
B_DL0 &= 50; \\
B_DL2 &= 0; \\
B_FL0 &= 32; \\
B_FL2 &= 23; \\
\text{% electric properties / measurement} \\
R_{ANE} &= 1.6; \quad \text{in Ohm} \\
R_{PHE} &= 1*0.2; \quad \text{in Ohm} \\
I_{amp} &= 0.1; \quad \text{in 10^8 A/cm^2} \\
freq &= 10; \quad \text{Hz} \\
R_{ANE} &= 3e-3; \quad \text{in volts} \\
\text{% time / simulation parameters} \\
sampling &= 100; \quad \text{% number of points per Period} \\
\text{numberOfPeriods} &= 10; \quad \text{% number of periods calculated for FFT} \\
\text{%%%%%%%%% setting up parameters %%%%%%%%%} \\
dt &= 1/(sampling * freq); \\
\text{time} &= (0 : dt : (sampling - 1) * dt)'; \\
\text{simTime} &= (0 : dt : numberOfPeriods / freq - dt)'; \quad \text{% total time} \\
\text{samplingFreq} &= (0 : (length(simTime) - 1)) * 1 / (dt * length (simTime))'; \\
I_{AC} &= I_{amp} * \sin(freq * simTime * 2 * \pi + \pi/2); \quad \text{% creation of one period}
\end{align*}
\]
theta_B = theta_B * pi / 180;
phi_B = phi_B * pi / 180;
[B_ext, scanLength] = createHysteresis(fieldSetArray, hysteresis);

periodLength = length(time);
equMagAngles = zeros(periodLength, 2);
V1f = zeros(scanLength, 1);
V2f = zeros(scanLength, 1);
V3f = zeros(scanLength, 1);

%%%%%%%%% simulation start %%%%%%%%%
for scanIterator = 1 : scanLength
  if B_ext(scanIterator) < 0
    initPhi = phi_B + pi;
  else
    initPhi = phi_B;
  end
  if scanIterator > scanLength / 2
    B_c = abs(B_c);
  else
    B_c = -abs(B_c);
  end
  if B_ext(scanIterator) <= B_c
    initTheta = pi - 0.1;
  else
    initTheta = 0.1;
  end
  options = optimoptions('fsolve', 'Display', 'off', 'Algorithm', 'trust-region-reflective');
  for i = 1 : periodLength %only first period
    B_DL0t = B_DL0 * I_AC(i);
    B_DL2t = B_DL2 * I_AC(i);
    B_FL0t = B_FL0 * I_AC(i);
    B_FL2t = B_FL2 * I_AC(i);
    [magAngles] = fsolve(@(magAngles) angleTorqueSum(magAngles, theta_B, phi_B, B_ext(scanIterator), B_k1, B_k2, B_DL0t, B_DL2t, B_FL0t, B_FL2t), [initTheta; initPhi], options);
    equMagAngles(i, :) = magAngles';
    initTheta = magAngles(1);
    initPhi = magAngles(2);
  end
  angles = repmat(equMagAngles, 10);
  V_H = I_AC .* (R_AHE .* cos(angles(:, 1)) + R_PHE .* sin(angles(:, 1))) + R_ANE .* cos(angles(:, 1)) - R_ANE .* sin(angles(:, 1)) .* cos(angles(:, 2));
  signalFFT = 2 * ifft(V_H);
  V1f(scanIterator) = abs(signalFFT(samplingFreq == freq)) * sign(real(signalFFT(samplingFreq == freq)));
  V2f(scanIterator) = -2 * abs(signalFFT(samplingFreq == 2 * freq)) * sign(real(signalFFT(samplingFreq == 2 * freq)));
end
V3f(scanIterator) = \(-1 \times \text{abs}(\text{signalFFT}(\text{samplingFreq} == 3 \times \text{freq})) \times \text{sign}(\text{real}(\text{signalFFT}(\text{samplingFreq} == 3 \times \text{freq})))\);

\begin{align*}
R1f &= \frac{V1f}{I_{amp}}; \\
R2f &= \frac{V2f}{I_{amp}}; \\
R3f &= \frac{V3f}{I_{amp}};
\end{align*}

if addNoise == true

\begin{align*}
\text{B}\_\text{ext}(i) &= \text{random('Normal', B}\_\text{ext}(i), \text{abs(B}\_\text{ext}(i)\times0.001), 1, 1); \\
\text{R1f}(i) &= \text{random('Normal', R1f(i), abs(R1f(i)\times0.001), 1, 1);} \\
\text{R2f}(i) &= \text{random('Normal', R2f(i), abs(R2f(i)\times0.001), 1, 1);} \\
\text{R}(i) &= \text{random('Normal', R3f(i), abs(R3f(i)\times0.001), 1, 1);} \\
\end{align*}

end

%%%%% Output creation %%%%%

\begin{align*}
\text{figure();} \\
\text{subplot(1, 2, 1);} \\
\text{plot(B}\_\text{ext}, \text{R1f},'o') \\
\text{xlabel('H_{ext} (Oe)')} \\
\text{ylabel('R -1f')} \\
\text{subplot(1, 2, 2);} \\
\text{plot(B}\_\text{ext}, \text{R2f},'o') \\
\text{xlabel('H_{ext} (Oe)')} \\
\text{ylabel('R -2f')} \\
\text{Output = table(B}\_\text{ext}, \text{R1f}, \text{R2f}, \text{R3f);} \\
\end{align*}

% % % createHysteresis % % %

\begin{align*}
\text{function } [\text{field}, \text{scanLength}] &= \text{createHysteresis( fieldSetArray, hysteresis)} \\
\text{field} &= []; \\
\text{for rowCounter} &= 1 : \text{size(fieldSetArray, 1)} \\
\text{fieldStart} &= \text{fieldSetArray(rowCounter, 1)}; \\
\text{fieldEnd} &= \text{fieldSetArray(rowCounter, 2)}; \\
\text{deltaField} &= \text{fieldSetArray(rowCounter, 3)}; \\
\text{fieldCounter} &= 0; \\
\text{if} &\text{ fieldStart > fieldEnd} \\
\text{deltaField} &= -\text{deltaField}; \\
\text{while} &\text{ fieldCounter \times deltaField + fieldStart > fieldEnd} \\
\text{field[end+1]} &= \text{fieldCounter \times deltaField + fieldStart}; \\
\text{fieldCounter} &= \text{fieldCounter + 1}; \\
\text{end} \\
\text{elseif fieldStart < fieldEnd} \\
\text{while} &\text{ fieldCounter \times deltaField + fieldStart < fieldEnd} \\
\text{field[end+1]} &= \text{fieldCounter \times deltaField + fieldStart}; \\
\text{fieldCounter} &= \text{fieldCounter + 1}; \\
\text{end} \\
\text{else} &\text{ error('your field Set Array has a stepsize of 0');}
\end{align*}
D.1. Harmonic Hall voltage simulation

end
field(end + 1) = fieldEnd;
end

field = field';
if hysteresis
    field = [field; flipud(field)];
end

i = size(field, 1);
while i >= 2
    if field(i) == field(i - 1)
        field(i) = [];
    end
    i = i - 1;
end
scanLength = length(field);

% angleTorqueSum

function T_sum = angleTorqueSum(magAngles, theta_B, phi_B, B_ext, B_k1, B_k2, B_DL0, B_DL2, B_FL0, B_FL2)

theta_M = magAngles(1);
phi_M = magAngles(2);
x = [1;0;0];
y = [0;1;0];
z = [0;0;1];
M = [sin(theta_M) * cos(phi_M); sin(theta_M) * sin(phi_M); cos(theta_M)];
B_ext_field = B_ext * [sin(theta_B) * cos(phi_B); sin(theta_B) * sin(phi_B); cos(theta_B)];

T_z = cross(M, B_ext_field); % Zeeman torque

% anisotropy
T_k1 = B_k1.*[cos(theta_M)*sin(theta_M)*sin(phi_M); -cos(theta_M)*sin(theta_M)*sin(phi_M); 0];
T_k2 = B_k2.*[cos(theta_M)*sin(theta_M)*sin(phi_M)^3; -cos(theta_M)*sin(theta_M)*sin(phi_M)^3; 0];
T_k = T_k1 + T_k2;

% current induced effects
T_DL = B_DL0 * cross(M, cross(y, M)) + B_DL2 * cross(z, M) * dot(M, x);
T_FL = (-B_FL0 - B_FL2 * dot(cross(z, M), cross(z, M))) * cross(y, M) - B_FL2 * cross(M, cross(z, M)) * dot(M, x);
T_sum = T_z + T_k + T_DL + T_FL;
end
D.2 Macrospin switching simulation

The following MATLAB code can be used to simulate SOT-induced magnetization switching in samples with PMA. It is based on the time integration of the LLG equation introduced in Section 2.3.1.

```matlab
%%%%%%%%% Parameter input %%%%%%%%%

% initial magnetization position
M(:, 1) = [0; 0; 1];

% external field stuff
theta_B = 90; % in degrees
phi_B = 0; % in degrees
B_ext = 0.1; % in T

% Magnetic properties
B_k = 0.6; % Anisotropy field in T
alpha = 0.5; % damping constant
gamma = 1.76*10^-11; % gyromagnetic ratio

% SOT's in T per 10^8 A/cm^2
B_DL0_amp = 0.05;
B_DL2_amp = 0.020;
B_FL0_amp = 0.032;
B_FL2_amp = 0.023;

% Current stuff
stabilizationTime = 0.2e-9; % Stabilization time before pulse injection in s
currentAmp = 10; % 10^8 A/cm^2
pulseLength = 0.5e-9; % duration of current pulse in s

% Simulation stuff
delta_t = 5e-13; % integration time step

%% simulation start %%

B_extField = B_ext * [ sind(theta_B)*cosd(phi_B);
                      sind(theta_B)*sind(phi_B);
                      cosd(theta_B)];

dmdt = 1;
inPulse(1) = false;
time(1) = 0;
itCounter = 1;
M = M / norm(M);

% going to equilibrium
while time(itCounter) <= stabilizationTime
    T_sum = vectorTorqueSum(M(:, itCounter), B_extField, B_k, 0, 0, 0, 0);
    [M(:, itCounter + 1), dmdt] = LLG_integration(M(:, itCounter),
                                              T_sum, delta_t, alpha, gamma);
    time(itCounter + 1) = time(itCounter) + delta_t;
    inPulse(itCounter + 1) = false;
    itCounter = itCounter + 1;
end
```
D.2. Macrospin switching simulation

```
pulseStartTime = time(itCounter);
B_DL0 = B_DL0_amp * currentAmp;
B_DL2 = B_DL2_amp * currentAmp;
B_FL0 = B_FL0_amp * currentAmp;
B_FL2 = B_FL2_amp * currentAmp;

% applying current pulse
while time(itCounter) <= pulseStartTime + pulseLength
    T_sum = vectorTorqueSum(M(:, itCounter), B_extField, B_k, B_DL0, B_DL2, B_FL0, B_FL2);
    [M(:, itCounter + 1), dmdt] = LLG_integration(M(:, itCounter), T_sum, delta_t, alpha, gamma);
    time(itCounter + 1) = time(itCounter) + delta_t;
    inPulse(itCounter + 1) = true;
    itCounter = itCounter + 1;
end

% going to equilibrium
eqCounter = 0;
while dmdt * delta_t >= 1e-15 && eqCounter < 5000
    T_sum = vectorTorqueSum(M(:, itCounter), B_ext_DMI, B_k, 0, 0, 0, 0);
    [M(:, itCounter + 1), dmdt] = LLG_integration(M(:, itCounter), T_sum, delta_t, alpha, gamma);
    time(itCounter + 1) = time(itCounter) + delta_t;
    inPulse(itCounter + 1) = false;
    itCounter = itCounter + 1;
    eqCounter = eqCounter + 1;
end

%%%%%%% output %%%%%%%
plot(time(:)',M(1,:))
hold on
plot(time(:)',M(2,:))
plot(time(:)',M(3,:))
xlabel('time [s]')
ylabel('M/ M_s')
legend('M_X','M_Y','M_Z','I-pulse')
data = [(time * 1e9)', M', inPulse'*currentAmp];

function T_sum = vectorTorqueSum(M, B_ext, B_k, B_DL0, B_DL2, B_FL0, B_FL2)
    x = [1; 0; 0];
    y = [0; 1; 0];
    z = [0; 0; 1];
    T_z = cross(M, B_ext); %Zeeman torque
    T_anis = cross(M, B_k * dot(M, z) * z); %anisotropy

%current induced effects
    T_DL = B_DL0 * cross(M, cross(y, M)) + B_DL2 * cross(z, M) * dot(M, x);
    T_FL = (-B_FL0 - B_FL2 * dot(cross(z, M), cross(z, M))) * cross(y, M) - B_FL2 * cross(M, cross(z, M)) * dot(M, x);```
D.2. Macrospin switching simulation

T_sum = T_z + T_anis + T_DL + T_FL;
end

function [M, dmdt] = LLG_integration(M, T_sum, delta_t, alpha, gamma)

dMdt = -abs(gamma) / (1 + alpha^2) * T_sum;
dMdt = dMdt - alpha * abs(gamma)/(1 + alpha^2) * cross(M, T_sum);
dmdt = norm(dMdt);
M = M + delta_t * dMdt;
M = M / norm(M);
end
E  OOMMF input file

Below we show a typical OOMMF input file in the MIF2.2 format used for the simulation of the static DW tilt shown in Fig. 5.2. Scalar and vector fields can be defined by color coding of a 24 bit bitmap image (see the OOMMF userguide for more information). We used this feature to define regions with different properties, e.g. magnetic and non-magnetic regions or out-of-plane and in-plane easy axes. Figure E.1 shows two typical bitmaps used herein. Red is the non-magnetic “universe”, white (black) is magnetic and initially saturated along $-z$ ($+z$).

The $MB_{-SOTevolve}$ extension allows to apply SOTs. It is based on the Oxs_EulerEvolve class and takes following additional input parameters. $J$ (default 0) is the current density in $10^8 \text{ A cm}^{-2}$. The current is always injected along $+x$ (this could be changed by defining it as an Oxs_vectorField object). $J_{\text{scaling}}$ (default 1) allows to inject spatially varying current densities ($J_{\text{injected}} = J \times J_{\text{scaling}}$). It is an Oxs_scalarField object and can be changed for each cell. $D_{L0}$ (default 50), $F_{L0}$ (default 32) and $F_{L2}$ (default 23) are the SOT amplitudes in mT per $10^8 \text{ A cm}^{-2}$ according to Eqs. (A.3) and (A.3). $\text{pulse\_start}$ (default 3.0e-10) defines the time in s at which the current pulse is injected, counting from the start of the simulation. $\text{pulse\_length}$ (default 2.0e-9) is the length of the rectangular current pulse given in s.

```
#MIF 2.2
#
# all units are SI

Parameter pulse_start 3.0e-10
Parameter pulse_length 2.0e-9
Parameter current_density 0.0
Parameter T_DL0 18.0
Parameter T_FLO 7.0
Parameter T_FL2 10.0
Parameter D 1.2
Parameter phi 90.0
Parameter theta 90.0
```

Figure E.1 | Bitmap image. Bitmap images used to define multiple regions in OOMMF.
Parameter B_ext 0.02

SetOptions [subst {
    basename test
    scalar_output_format %.17g
    vector_field_output_format {text %.6g}
}]

set PI [expr {4*atan(1.1)}]
set MU0 [expr {4*$PI*1e-7}]

Specify Oxs_ImageAtlas:sampleAtlas {
    xrange {0 1500e-9}
    yrange {0 1500e-9}
    zrange {0 1e-9}
    viewplane xy
    image anisotropic_CIDWM_v2.0.bmp
    colormap {
        red universe
        black up
        white down
    }
}

set cellsize 1e-9

Specify Oxs_RectangularMesh:mesh [subst {
    cellsizes {[expr 4*$cellsize] [expr 4*$cellsize] $cellsize} [subst {
        atlas :sampleAtlas
    }}]

Specify Oxs_UniaxialAnisotropy {
    K1 { Oxs_AtlasScalarField {
        atlas :sampleAtlas
        default_value 0.0
        values {
            up 657.0e3
            down 657.0e3
        }
    }]
    axis { Oxs_AtlasVectorField {
        atlas :sampleAtlas
        default_value {0 0 1}
        values {
            up {0 0 1}
            down {0 0 1}
        }
    }]
}

Specify Oxs_UniformExchange {
    A 1e-11
}

set DD [expr {$D/1000}]

Specify Oxs_DMExchange6Ngbr:DMEx [subst {
    default_D 0.0
    atlas :sampleAtlas
    D {
        up up $DD
    }
}]
down down $DD$

\[
\begin{align*}
\text{set } & B_x \left[ \text{expr} \left\{ \frac{B_{\text{ext}} \times \cos(\phi/180 \times \pi) \times \sin(\theta/180 \times \pi)}{} \right\} \right] \\
\text{set } & B_y \left[ \text{expr} \left\{ \frac{B_{\text{ext}} \times \sin(\phi/180 \times \pi) \times \sin(\theta/180 \times \pi)}{} \right\} \right] \\
\text{set } & B_z \left[ \text{expr} \left\{ \frac{B_{\text{ext}} \times \cos(\theta/180 \times \pi)}{} \right\} \right]
\end{align*}
\]

Specify Oxs_FixedZeeman:field [\text{subst} \{\text{multiplier} \left[ \text{expr} \left\{ \frac{1}{\mu_0} \right\} \right] \text{field} \{B_x \ B_y \ B_z\} \}]

Specify Oxs_Demag {};

Specify MB_SOTevolve [\text{subst} \{\alpha 0.5 \gamma_G 2.211e5 \text{do_precess 1} \text{min_timestep 1e-15} \text{max_timestep 1e-10} \text{start_dm 0.01} \text{J \$current\_density} \text{pulse\_start \$pulse\_start} \text{pulse\_length \$pulse\_length} \text{DL_0 \$T\_DL0} \text{FL_0 \$T\_FL0} \text{FL_2 \$T\_FL2} \}]

Specify Oxs_TimeDriver [\text{subst} \{\text{evolver MB_SOTevolve} \text{mesh :mesh} \text{stage\_count 150} \text{stopping\_time 20e-12} \text{Ms \{ Oxs_AtlasScalarField \{atlas :sampleAtlas default\_value 0 values \{up 900e3 down 900e3\} \}\} \text{m0 \{ Oxs_AtlasVectorField \{atlas :sampleAtlas default\_value \{0 0 1\} values \{up \{0 0 1\} down \{0 0 -1\}\} \}\} \} \}]

}]}
Bibliography


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Bibliography


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