A new route towards synchronous domain wall motion
studying precession torque effects in nanostrips with perpendicular magnetic anisotropy

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A new route towards synchronous domain wall motion

STUDYING PRECESSION TORQUE EFFECTS IN NANOSTRIPS WITH PERPENDICULAR MAGNETIC ANISOTROPY

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Abstract

Magnetic domain wall motion is an intensively studied subject because of its possible applications for data storage. High density, non-volatile computer memories could be fabricated if novel concepts based on domain wall motion are realized. In this thesis domain wall motion, using an unconventional driving mechanism based on the precession torque, is studied.

Precession torque driven domain wall motion in magnetic nanostrips with perpendicular magnetic anisotropy (PMA), is investigated both theoretically and experimentally. The theoretical investigation is done using a collective coordinates model and object oriented micromagnetic framework simulations. Both methods show that when an in-plane magnetic field pulse is applied to a domain wall in a PMA system, the domain wall is indeed moved by the precession torque. When realistic values of the material parameters are used, high domain wall velocities, in the order of 100 m/s, are predicted. The influence of the strength of the applied pulse, Gilbert damping parameter and the chirality of the domain wall are investigated and the observed behaviour is discussed.

Experimentally, we study the depinning of domain walls as a function of the applied in-plane field pulse and the results are compared to quantitative predictions from the model. The observed behaviour does not correspond to the predictions. Further analysis indicates that in order to observe effects from the precession torque, pulses with a rise time in the order of a nanosecond are required, a requirement not fulfilled in our setup.

During the experiments, we observed a chiral dependence of the domain wall depinning on static in-plane fields along the nanostrip. This unexpected behaviour can be explained by the Dzyaloshinskii-Moriya interaction, DMI. This allows for a novel way to study the DMI, using well-defined structures and relatively simple measurements.
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One of the most famous expressions in the computer industry is “640 KB ought to be enough for anybody”, a quote from 1981 by no one less than Bill Gates. Though up to present day it remains unclear whether Bill Gates actually said these words, the quote illustrates very well how rapidly the use of data has grown, and how difficult it was to imagine this on before hand. Computers nowadays are hardly comparable to the ones in 1981; figure 1.1 shows an “Apple III” an example of a computer released in that year, with a maximum memory of 512 KB. Nowadays, an ordinary smart phone has orders of magnitude more memory, for instance the “Samsung Galaxy Express”, also shown in figure 1.1, has 8 GB available for storage. The expectation is that the demand for more data per volume, and faster access to this data, will keep increasing. One expectations is the development of the internet of things [1]; this refers to a network of a vast amount of every day objects which generate data and communicate this. This is only one example of a development that would be accompanied by a increase in usage of data. This means research on techniques for data storage will have to keep up with the pace of this growing demand.

One research area that often finds applications in the data storage industry is spintronics. It focusses on devices based not only on the charge of electrons, but also on their spin, an intrinsic form of angular momentum. Just as the angular momentum arising from an orbital motion of a charged particle, the spin has an associated magnetic moment. It was realized that involving magnetic properties in electronics could lead to improvements; higher data densities and non-volatile memories could become possible. This research field had a breakthrough in the 1980’s, when giant magnetoresistance (GMR) was discovered [2]. This discovery that was implemented in industry only a few

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Over the years many new techniques have been developed to increase the density and rate of data, and we would like to highlight one; the racetrack memory, invented by Stuart Parkin [3]. On a conventional hard disk drive, data is stored in magnetic bits, where a magnetization in the one direction represents a “0” while in the other direction it represents a “1”. The state of such a bit can be read out by a read head based, for instance based on GMR. To read out a certain bit somewhere on the hard disc, the read head has to be exactly above this bit. To accomplish this, the hard disc has to rotate and the read head will have to translate, a process which takes time and causes heating of the computer. For the writing of a bit the same problem occurs. In the racetrack memory, the information is stored in a magnetic wire, with different regions, called domains, where the magnetization can be in different directions. If these magnetic domains can be moved through the wire, to the position of the fixed read head, a solid state computer memory without any moving parts can be realized. Another advantage is that the wires could be folded in to the third dimension, making a very dense storage of data possible.

The key challenge for this new memory is being able to controllably move the magnetic domains through the wires or, equivalently, move the boundaries between the domains, the domain walls. Therefore the idea of a racetrack memory has greatly stimulated the research on domain wall motion. Domain wall motion by sending a current through the wire was extensively studied by various research groups [4]. This has let to the production of systems that could indeed function as a racetrack memory, but also to new physics; unexpected results were obtained and explained by exotic new phenomena like the spin Hall effect, the Rashba effect and the Dzyaloshinskii-Moriya interaction. The debate about these effects and their importance in current-driven domain wall motion continues until the present day. Another approach was altering the properties of the
magnetic wire in such a way that it becomes a ratchet for domain walls, which makes a racetrack memory driven by magnetic fields possible [5]. Within this lively research field, we study another, new, approach for domain wall motion, using short pulses of magnetic fields to exert a precession torque on the domain wall. This way, we hope to combine advantages of earlier approaches and, of course, get a deeper understanding of the physics involved.

The goal of this project is the investigation of precession torque driven domain wall motion, and to present the work in a clear way this thesis is divided into chapters. A theoretical background will be given in chapter 2. This includes an introduction to the relevant energy terms, types of domain walls are discussed and conventional field and current driven domain wall motion are discussed. Further an intuitive description of precession torque driven domain wall motion is given, after which a model to investigate this phenomenon is build. This is followed by a more practical part, chapter 3, in which the detailed design of our samples is presented as well as the production procedure. After this, our measurement routine is presented together with an introduction to the crucial tool in our measurements; the Kerr microscope. In chapter 4 results that are obtained using our model and computer simulations are presented and discussed. First, the general behaviour and the influence of varying relevant parameters is studied. Also the possibility to move various types of domain walls using the precession torque is discussed. The model is adapted to closely match our experimental routine, and quantitative expectations for the measurements are obtained. Last, precession torque driven domain wall motion turns out to have one unwanted attribute, and solutions for this are discussed. In chapter 5 our experimental results are presented in a chronological fashion. Extra attention is given to unexpected results, which turn out to be related to an exiting phenomenon called Dzyaloshinskii-Moriya interaction. Finally, we summarize our conclusion in chapter 6 and give an outlook on future possibilities involving precession torque driven domain wall motion in chapter 7.
Chapter 2

Theory

2.1 A new way of domain wall motion

2.1.1 Introduction to domain walls

Magnetic materials usually consist of multiple areas with a different magnetization, called magnetic domains. On the boundaries of the areas, the magnetization varies continuously. The region where this happens is called the domain wall. In this work, domain walls in nanostrips with perpendicular magnetic anisotropy (PMA), which means that inside the domains the magnetization points out of the plane, are studied. The width and length of the wires are in the order of micrometers and the thickness in the order of one nanometer. A schematic representation of such a magnetic nanostrip with a domain wall is shown in figure 2.1 (a). In this figure, a coordinate system is defined as well; the $z$ direction is perpendicular to the plane of the nanowire and the $x$ direction is along the nanostrip. These definitions will be used throughout this thesis.

Figure 2.1: (a) Schematic view of a part of a PMA system with a domain wall and the definition of axes that will be used throughout this report. (b) Schematic representation of a racetrack memory, figure adapted from literature [3].
Our goal

Research on domain walls is often aimed at applications in the data storage industry. Here we discuss a novel type of memory device, which inspired the research in this thesis. This device is the *racetrack memory*, already briefly mentioned in the introduction, which is schematically shown in figure 2.1 (b). A nanowire consisting of blue and red parts is shown. These colours represent domains in which the magnetization points in a certain direction. In the bottom part of the figure, two small constructions attached to the nanowire are shown, which represent a reading and writing device. Here the direction of the magnetization can be read out and interpreted as a “0” or a “1”, or the direction of the magnetization can be changed. Mechanisms to do this are for instance magnetic tunnel junctions for sensing [6], and spin transfer torque for writing [7], but the reading and writing will not be the focus point of this report.

One advantage of this new device can be seen in figure 2.1; the nanostrip does not have to lay flat on the substrate but can be folded into the third dimension. This could greatly increase the amount of data per surface area, but the production of such a folded nanowire remains a technological challenge until the present day, which will not be addressed in this thesis. The other, even more innovative, advantage is that a solid state computer memory without any mechanically moving parts can be realized! The reading and writing heads are fixed and can only access the domain which is directly above them. The key idea of the racetrack memory is that the magnetic domains can be shifted through the wires; in figure 2.1 (b) the domain structure in the right picture is shifted with respect to the left picture. This also mean that an other bit is accessible for the read and write head. This has potentially huge advantages with respect to the heat generation, fragility and speed at which data can be accessed. For the realisation of such a memory one should be able to controllably move the domains, or equivalently the domain walls, through the nanostrip. In the original paper it is proposed to do this by sending a spin polarized current through the nanowire. In this report, an other approach, relying on a completely different physical mechanism, is investigated.

Energies in our magnetic nanostrip

We now take a closer look at the nanostrips to describe the magnetism in and identify the energy terms that are of importance. Crucial for the magnetic properties of magnetic materials are the electrons and the magnetic moment associated with their intrinsic angular momentum called *spin*. Electrons occupy certain orbitals around a nucleus. They usually do this in pairs, and the two electrons of a pair have an opposite spin, making the total magnetic moment zero. Ferromagnetic elements have one or more
unpaired electrons, making a net magnetic moment possible. In our nanostrips, the ferromagnetic material is an ultra thin (0.6 nm) layer of cobalt, sandwiched between two platinum layers. Magnetic domains, like in a racetrack memory, form when spins of neighbouring atoms align. The mechanism causing this is the exchange interaction; a consequence of a combination of Coulomb interaction and the Pauli exclusion principle. The result of this interaction is schematically shown in figure 2.2 (a), and a formal description of this energy is given by equation 2.1.

\[ E_{\text{exchange}} = \frac{A}{M_s^2} \left( \nabla \vec{M} \right)^2 \]  

(2.1)

**Figure 2.2:** Schematic cross section of our nanostrip indicating the direction of the magnetic moments due to the (a) exchange interaction (b) external applied magnetic field (c) the demagnetization field (d) magnetic anisotropy.

In this equation, \( A \) is the exchange stiffness, \( \vec{M} \) is the magnetization and \( M_s \) is the saturation magnetization which is the length of the magnetization vector and is a material constant. This term is minimal when there are no gradients in the magnetization. An additional energy contribution arises when an external magnetic field \( \vec{H}_{\text{ext}} \), is applied, which is called Zeeman energy. The energy is minimized when the magnetization aligns with the applied field, as described by equation 2.2:

\[ E_Z = -\mu_0 \vec{M} \cdot \vec{H}_{\text{ext}} \]  

(2.2)

Here \( \mu_0 \) is the permittivity in vacuum and \( \vec{M} \) again the magnetization.

The exchange interaction is isotropic; it makes neighbouring spins align, but there is no preference for any direction. However, as mentioned at the start of this chapter and shown in figure 2.1, the magnetization in our nanostrips prefers the direction perpendicular to the plane of the nanostrip. This preference for a certain direction is called magnetic anisotropy. On first sight perpendicular magnetic anisotropy is not expected for ultrathin films, considering the demagnetization fields. This is a field generated by the magnetization itself: \( \vec{H}_d = -N \vec{M} \). Intuitively a magnetic body (consisting of only on domain), is build up from tiny magnetic moments, all generating dipole fields. Therefore, at the surface of the body, “free” magnetic poles can exist, generating this demagnetization field. The \( N \) in the formula is the demagnetization factor, which is determined by the shape of the magnetic body, and the direction of the magnetization. A
formula for $N$ for arbitrary shaped bodies does not exist, but for the simple geometry of a thin film it can be calculated using Ampere’s and Gauss’s laws, as is done in literature [8]. This results in $N = 0$ when the magnetization is in the plane (the poles are infinitely far apart) and $N = 1$ when the magnetization is out-of-plane. Just as with externally applied fields, the energy is minimized when the magnetization aligns with the field. But the magnetization is antiparallel to the demagnetization field per definition, so it is energetically favourable if $N$ is as small as possible. Following this argument, a thin film should always have an in-plane magnetization, like in figure 2.2 (c).

This argument neglects contributions from surfaces or interfaces. Néel realized this already in 1954 and predicted that for systems with small dimensions these contributions could be of importance [9]. When it became possible to experimentally produce ultrathin structures, for example using sputtering which will be discussed in chapter 3, systems with perpendicular magnetic anisotropy were produced indeed [10]. The mechanism that makes PMA possible in certain ultrathin films is spin-orbit coupling. An intuitive interpretation of this effect would be the following; the electron orbits the nucleus, and this nucleus creates an electric field because it has a charge. To describe the fields as observed from the rest frame of the electron, a Lorentz transformation has to be performed; this shows that the electron also experiences a magnetic field next to the electric field. The spin of an electron is sensitive to these magnetic fields according to the Zeeman energy. This is why a coupling exists between the spin and orbit of an electron. The orbits of the electrons are affected by the surroundings of the atoms, here the potential created by the other atoms in the lattice. The surroundings are different at an interface and in the bulk, and it is possible that at the interface orbitals exist that are quenched in the bulk. At certain interfaces, between a magnetic material and a material with large spin-orbit coupling like platinum or palladium, this result more favourable out-of-plane orientation of the magnetization, see figure 2.2 (d). Because this is an interface effect it becomes increasingly important for thinner films. This was a very brief explanation for PMA in ultrathin films. More detailed descriptions exist [11], but are not required to understand the physics in this thesis. In our description, both the interface contribution of spin-orbit coupling and the demagnetizing energy are combined in an effective uniaxial anisotropy, leading to an energy contribution given in equation 2.3.

$$E_K = K_{eff} \sin^2 \theta$$  \hspace{1cm} (2.3)

$K_{eff}$ is the effective anisotropy constant and $\theta$ is the angle between the magnetization and the easy axis, in our case the $z$ axis. Note that more possible contributions to the anisotropy exist. crystalline anisotropy in the bulk, exchange bias due to coupling to
an antiferromagnetic layer, strain induced anisotropy or anisotropy induced by growth in an external magnetic field. These contributions are not present or not dominant in the samples studied in this thesis [12].

So why do we focus on PMA systems if they only occur under special conditions? The motivation for this lies with the application for computer memories. In order to be interesting for industry, the scalability is of importance. When data is stored as magnetic domains, a limiting quantity when scaling down is the size of the region separating two domains, the domain wall. The size of a domain wall is determined by the exchange interaction and anisotropy, both introduced earlier in this section. Exchange interaction favours wide domain walls, because this results in small misalignments for neighbouring spins. Anisotropy favours small domain walls, because this reduces the number of spins that is not aligned with the easy axis. Minimizing the domain wall energy results in equation 2.4:

\[
\lambda \sim \sqrt{\frac{A}{K_{\text{eff}}}}
\]  

(2.4)

Here \( \lambda \) is the domain wall width. The energy of the domain wall when it takes on this width is given in equation 2.5.

\[
E_{\text{DW}} = 4\sqrt{AK_{\text{eff}}}
\]  

(2.5)

For in-plane materials the width of a domain wall is in the order of 100 nm while for PMA materials it is around 10 nm.

The discussion in this subsection concerned a magnetic nanostrip in the static case and a domain wall at a fixed position, but next we will introduce how domain walls can be moved.

**Conventional domain wall motion**

In general there are two ways to do move domain walls; using magnetic fields or using electrical currents. The first possibility for moving domain walls is applying a magnetic field along the easy axis of the magnetic material, so for PMA materials along the \( z \) direction. The driving energy for this is the Zeeman energy, defined in equation 2.2. This will make a domain wall propagate in such a way that a domain with the magnetization in the direction of the applied field increases in size. On the contrary, domains with the magnetization anti-parallel to the field will shrink. After a while the magnetization will
point in the direction of the applied field everywhere in the structure. For a racetrack memory this would mean all information is lost, highly undesirable.

Two regimes can be distinguished with respect to the dependence of the domain wall velocity on the applied field, called the creep and flow regime. Weak disorder in a system will exert a pinning force on the domain wall. In the flow regime, the driving force of the magnetic field is much larger than this pinning force, and the disorder does not affect the domain wall motion. The domain wall velocity is proportional to the applied field. If the driving force is smaller than the pinning force, the domain wall motion is thermally activated, which is called creep. The velocity is now described by the creep law.

\[ V = V_0 \exp \left( -\alpha H_z^{-\mu} \right) \]  

(2.6)

This law describes how the velocity \( V \), of a domain wall depends on the applied out-of-plane field \( H_z \), where \( V_0 \) is the characteristic speed, the exponent \( \mu \) is typically 1/4 and \( \alpha \) is a scaling constant. Though this method of domain wall motion is not the focus point of this thesis, it plays a key role in our experimental routine, as will become clear later. The flow regime will start above applied field of 60 mT for the materials we use [13], while for our experiment fields in the order of 10 mT are used. We therefore expect to be in the creep regime, but because we cannot measure the domain wall velocity in our setup, this cannot be checked.

Another way of inducing domain wall motion in a magnetic nanowire is sending a spin polarized current through it, which is the approach used in the original racetrack paper. An intuitive explanation for this is the following: the spins of the conduction electrons follow the magnetization of the material they flow through. This means that they change from up to down (or the other way around) when passing a domain wall. Conservation of angular momentum dictates that a torque will be exerted at the localized spins in the domain wall, which will make the domain wall move in the direction of the electron flow. This is called the adiabatic spin-transfer torque [14].

When experiments were conducted, it turned out this picture of current driven domain wall motion is oversimplified. Domain wall motion with large velocities or in the direction opposite to the electron flow were observed [4]. Various effects have been introduced in order to explain the observed results; the non-adiabatic spin transfer torque [15], the Rashba effect [16] and the spin Hall effect [17]. The discussion about the importance of these effects for current-driven domain wall motion continues to the present day. In spite of these unexpected complex dynamics, synchronous motion of multiple domain walls, the requirement for a racetrack memory, could be achieved. A disadvantage of
2.1 A new way of domain wall motion

this way of domain wall motion is that relatively high current densities through the nanostructures are necessary, causing significant heating [18].

2.1.2 Precession torque driven domain wall motion

The two main methods to drive domain walls both have an important drawback: current-induced motion requires large currents through the nanostructures, and field driven motion eventually erases an alternating domain structure. In this thesis, an alternative driving force for domain wall motion is presented, combining the advantages of both aforementioned approaches. Before a mathematical model is presented, a intuitive explanation of the basic principle is given.

In conventional field driven domain wall motion the magnetic field is applied along the easy axis, and the wall moves to reduce the Zeeman energy, as explained in section 2.1.1. In our PMA systems, the easy axis is the $z$ axis. But what will happen when the magnetic field is applied along the $x$ or $y$ axis? From equation 2.2 it can be seen that the Zeeman energy does not favour up or down domains depending on a field applied in the $x$ direction, so at first glance this does not drive domain wall motion.

However, it can be seen that in-plane fields do affect the system, when examining the Landau-Lifschitz-Gilbert (LLG) equation. This is a widely used dynamical equation for the magnetization developed by Landau and Lifschitz in 1935, and adapted by Gilbert in 1954 to describe the damping in a more logical way. The LLG equation is given in equation 2.7, without any current induced terms because these are not relevant in this work.

$$\frac{\partial \vec{M}}{\partial t} = -\gamma \vec{M} \times (\vec{H}_{ext} + \vec{H}_{eff}) + \alpha \frac{\vec{M}}{M_s} \times \frac{d\vec{M}}{dt} \quad (2.7)$$

In this equation $t$ represents the time, $\gamma$ the gyromagnetic ratio, $\alpha$ the Gilbert damping parameter, $\vec{H}_{ext}$ the externally applied magnetic field and $\vec{H}_{eff}$ is an effective field the spins experience, accounting for the exchange interaction and anisotropy.

The first term, which is perpendicular to both the magnetic field and the magnetization, describes the precession around the applied field. The second term, perpendicular to the magnetization and its time derivative see figure, is a phenomenological damping which makes the magnetization align with the field. In figure 2.3 (a) the effect of both terms is presented in a cartoon. The precession term is the key term that makes domain wall motion by in-plane fields possible, and can be derived using basic quantum mechanics.
The hamiltonian $\hat{H}$ of a spin in an external magnetic field is given in equation 2.8, which is equivalent to the Zeeman energy.

$$\hat{H} = \gamma \mu_0 \hat{\mathbf{S}} \cdot \mathbf{H}_{\text{ext}}$$  \hspace{1cm} (2.8)

$\hat{\mathbf{S}}$ represents the expectation value of the spin operator. Note that the hats in this equation refer to operators, and not to unit vectors like elsewhere in this thesis. An expression for the time derivative of the spin operator can be obtained from the Ehrenfest theorem. Together with the fundamental commutation relations for spin, this leads indeed to the first term of the LLG equation (after replacing $\hat{\mathbf{S}}$ by $\hat{\mathbf{M}}$), see equation 2.9.

$$\frac{d\hat{\mathbf{S}}}{dt} = \frac{1}{i\hbar} \left[ \hat{\mathbf{S}}, \hat{\mathbf{H}} \right] = -\gamma \mu_0 \hat{\mathbf{S}} \times \mathbf{H}_{\text{ext}}$$  \hspace{1cm} (2.9)

The effect of an in-plane field can be understood considering the configuration in figure 2.3 (b), which shows the top view of a PMA nanowire with a domain wall. When a magnetic field is applied in the $x$ direction, the LLG equation can be applied on the magnetic moments inside the wall; they now point along the $y$ direction, so according to equation 2.7 they will rotate around this field and at a certain moment point in the positive $z$ direction. However, when the magnetic moments point along the $z$ direction, they are no longer part of the domain wall but have become part of the blue domain; the domain wall has moved! This argument is only intuitive, and it raises some questions. What happens when the magnetic field is turned off? Why are the spins outside the domain wall not affected by a torque due to the in plane field? How does this make synchronous domain wall motion possible, if the direction of the spins in the domain wall is random? To address these questions and to support this intuitive idea, a theoretical model and simulation program will be introduced in the next subsections.

**Figure 2.3:** (a) Precession of a magnetic moment inside a magnetic field. (b) In the configuration of the nanowire depicted, the spins in the domain wall will experience this torque and start to precess, resulting in domain wall motion.
2.1 A new way of domain wall motion

2.1.3 The 1D model

Up to now only an intuitive argument explaining precession torque driven domain wall motion has been given. In this section a more formal model is developed to investigate this phenomenon. For this, the so-called collective coordinates approach is used. We build a model in which our system is described as an infinite wire in which a magnetic domain wall is present. We assume that the thickness and width of the wire are small, so the magnetization is constant in these directions, and we have a one dimensional system. Also, we assume that the domain wall is a rigid body, it cannot get wider or narrower and its “shape” is fixed. The shape can mathematically be defined as the out of plane angle of magnetization, $\theta$, as a function of the position. In this model, the domain wall is described by only two parameters; the position along the $x$ axis, $q$, and the tilting angle of the magnetization out of the easy plane at the center of the domain wall, $\phi$. A schematic representation of the situation is shown in figure 2.4. We first calculate the shape of the domain wall in the static case, and then assume this shape is not altered in the dynamic case. This calculation is given in detail in Appendix A.1. The result is given in equation 2.10.

\[
\cos \theta = \tanh \left( \frac{p(x - q)}{\lambda} \right)
\]  

(2.10)

Here, $\lambda$ is the domain wall width, introduced in equation 2.4. Further, $p$ accounts for the polarity of the domains and can be +1 for the polarity we assume, and -1 to describe the opposite polarity. To calculate what happens when external fields are applied, the method developed by Thiele is used [19]. Every term in the LLG equation can be written in the form $-\gamma \vec{M} \times \vec{H}$. Now the sum of all these “fields”, $\vec{H}_{tot}$, should be equal to 0. We can identify a force, $f$, and torque, $\tau$ as $f = -\vec{H}_{tot} \cdot (\frac{\partial \vec{M}}{\partial x})$ and $\tau = -\vec{H}_{tot} \cdot (\frac{\partial \vec{M}}{\partial \phi})$. Because $\vec{H}_{tot} = 0$, this force and torque are equal to 0 as well, which can be used to derive equations for $q$ and $\phi$. Detailed calculations can be found in Appendix A.2, but here only the resulting equations of motion are presented.
\[ \alpha \dot{q} - p \lambda \dot{\phi} = 0 \]
\[ p \dot{q} + \alpha \lambda \dot{\phi} = \frac{\pi \gamma \lambda}{2} (H_x \cos \phi + H_y \sin \phi) - \frac{\gamma \lambda H_D}{2} \sin 2\phi \]  (2.11)

This can be rewritten to
\[ \begin{pmatrix} \dot{q} \\ \dot{\phi} \end{pmatrix} = \frac{1}{1 + \alpha^2} \begin{pmatrix} p \lambda \\ \alpha \end{pmatrix} \left( \frac{\pi \gamma}{2} (H_x \cos \phi + H_y \sin \phi) - \frac{\gamma}{2} H_D \sin 2\phi \right) \]  (2.12)

\( H_x \) and \( H_y \) are the external fields applied in the \( x \) and \( y \) direction, respectively. Which of these is needed to move the domain wall depends on the direction of the spins inside the domain wall. \( H_D \) is a field resulting from demagnetization effects, the sign of which determines what in-plane angle is favorable for the spins inside the wall. These equations do not have an analytical solution for \( q \) and \( \phi \) as a function of \( t \). Numerical solutions are possible, and the physics following from this will be presented in the section 4.1.

### 2.1.4 Object Oriented Micromagnetic Framework

Another way in which the system can be investigated is by simulation. The simulation program that is used is called OOMMF which is short for the object oriented micromagnetic framework [20]. This freely available software, developed by NIST, is widely used for micromagnetic studies. At the time this thesis was written, 1597 papers cited OOMMF. A few examples of subjects that are studied using OOMMF are current driven domain wall motion [21], novel spin structures [22] and applications like biosensors for the detection of magnetic beads [23]. OOMMF solves the LLG equation numerically. It is finite difference simulation, which means that the investigated system subdivided into small cuboids (we use a cellsize of 1 nm\(^3\)). In literature ([17], supplementary information) realistic material parameters corresponding to our nanostrips can be found: the anisotropy constant \( K = 1.5 \times 10^{6} \) J/m\(^3\), the exchange stiffness \( A = 1.6 \times 10^{-11} \) J/m and the saturation magnetization \( M_S = 1.4 \times 10^{6} \) A/m.

It is interesting to investigate our system using both OOMMF and the 1D model because for both methods other approximations are made. In the 1D model it is assumed that the magnetization does not vary along the width of the wire and the shape of the domain wall is fixed, OOMMF does not have these restrictions. Also effect of the in-plane field on the spins outside the domain wall can be observed by OOMMF. Other differences are that in the 1D model the wire is truly infinite, while in OOMMF it is not. This means that edges effects, for instance an additional demagnetization energy, can occur in the
2.1 A new way of domain wall motion

results calculated by OOMMF. A more practical difference is the following: some of the OOMMF simulations presented in this thesis, have to run for more than 24 hours, so these simulations are not convenient for investigating behaviour on longer timescales. In the collective coordinate model calculations however, the behaviour of the domain wall up to microseconds after applying the in-plane field can be obtained within one minute. In section 4.1 results obtained from OOMMF simulations will be presented, and compared to predictions from the model.

2.1.5 Types of domain walls

Different types of domain walls can be identified based on the orientation of the spins inside the wall. While deriving the 1D model, no attention was paid to this point; the out-of-plane angle, \( \theta \) was calculated as a function of the position, but the most favorable in-plane angle, \( \phi \) has not been calculated. We assumed this to be 0 rad, so the spins in the domain wall point along the \( y \) axis. This spin configuration is called a Bloch wall, see figure 2.5 (a). Another type is a wall with spins in the \( x \) direction, which is called a Néel wall, see figure 2.5 (b). Which configuration is most favorable depends on the geometry of the wire. In a simplified model, the domain wall can be approximated as a bar with one single spin. In a bar, the demagnetization field will make this spin point along the longest axis. This means that nanowires much wider than the domain wall width (say more than 50 nm) will have Bloch walls, very narrow nanowires will have Néel walls. This idea has been confirmed by experimental observations in literature [24].

![Figure 2.5: Top view nanowires with (a) a Bloch wall (b) a Néel wall.](image)

We also have to consider a second property in which domain walls can differ; in figure 2.6 the topviews of four PMA nanowires, all with Bloch walls, are depicted. They differ with respect to the in-plane angles of the domain walls and the polarity of the domains. The four configurations can be divided into two groups based on the way the spins rotate. We say that the two groups have a different chirality, and call one group left-handed and the other one right-handed. Using the intuitive interpretation of precessional motion described in subsection 2.1.2, it can be derived that when a magnetic field is applied in the \( +x \) direction, domain walls with a left-handed chirality will move to the right, and domain walls with a right-handed chirality will move to the left. So when multiple domain walls are present in one wire, synchronous motion is only possible if the chirality is the same for all domain walls. How this can be achieved will be addressed later in this report.
2.1.6 Switching

Up to now numerous considerations involving domain walls have been presented. But how a domain wall can be introduced in a system, has not been discussed. When we have a uniform magnetic body, a domain wall is created when in a part of the body the magnetization reverses or, in other words, switches. The simplest model to describe switching is the Stoner-Wohlfarth model, in which the whole magnetic body per definition consists of a single domain. When this single domain is placed in a magnetic field, its energy of is the sum of the Zeeman energy and anisotropy. If the body has an perpendicular magnetic anisotropy, like our samples, the energy when magnetic fields are applied in the out-of-plane direction is:

\[ E = -\mu_0 M_S H_z \cos \theta + K_{eff} \sin^2 \theta \]  

(2.13)

This equation is obtained by combining equation 2.2 and 2.3. It can be seen from this equation, that the energy is minimal when the magnetization is aligned with the applied field, corresponding to \( \theta = 0 \) or \( \theta = \pi \). However, when a very small magnetic field changes from positive to negative, the magnetization will not follow. This is because the energy barrier caused by the anisotropy. To calculate the the moment when the magnetization switches derivative of equation 2.13 has to be computed. When this value becomes zero for a certain angle \( \theta \), there is no longer a local minimum for this angle, and the magnetization switches. The field at which this happens, according to this argument, is derived in equation 2.14.
2.1 A new way of domain wall motion

\[ \frac{\partial^2 E}{\partial \theta^2} = \mu_0 M_S H_z \cos \theta + 2K_{eff} \left( \cos^2 \theta - \sin^2 \theta \right) \]
\[ \left. \frac{\partial^2 E}{\partial \theta^2} \right|_{\theta=0, \pi} = \pm \mu_0 M_S H_z + 2K \]
\[ H_{\text{switch}} = \pm \frac{2K}{\mu_0 M_S} \]

Now a theoretical prediction of the magnetization as a function of the applied field can be computed, and is shown in figure 2.7. This is known as a hysteresis loop, and such curves are often measured in the experimental part of this project. It can be seen that the magnetization depends on the history of the system, sweeping the applied field from positive to negative values gives a different curve than sweeping the applied field from negative to positive values. The arrows on the curve indicate in which direction the applied field is swept. A jump in the magnetization occurs at the switching field, and this feature is mainly used in the experimental investigation.

![Figure 2.7: Hysteresis loop as calculated using the Stoner-Wohlfarth model.](image)

When the switching field for our samples is calculated using equation 2.14, a value of about 2 Tesla is found, two orders of magnitude higher than what is experimentally observed. This is an indication that the situation is more complex and the assumptions made in the Stoner-Wohlfarth model are not justified. The magnetization of a body usually does not behave like a single spin (with the exception of systems with dimensions smaller than the domain wall width in all directions). The anisotropy has local deviations because of imperfections of the magnetic structure. The magnetization usually switches in a small area at a point where the anisotropy is low, which is called nucleation. At what
field strength nucleation occurs can be calculated as well; the inverted domain lowers the Zeeman energy, but at the same time a domain wall is created. The latter results in an increase of energy, due to the anisotropy and exchange interaction. Nucleation can only occur when the profit in Zeeman energy outweighs the costs in exchange energy, which is formally given in equation 2.15:

\[ E/t = 2\pi R\sigma_0 - 2\mu_0 M_S H_z \pi R^2 \]  

(2.15)

Here \( R \) is the radius of the reversed domain and \( \sigma_0 \) is the domain wall energy density. The radius for which this energy is equal to zero, is the minimal radius for which the reversed domain is stable. The energy barrier for the formation of a reversed domain is equal to the energy cost of a domain wall around a domain with this radius:

\[ E_{\text{barrier}} = \frac{\pi \sigma_0^2 t}{2\mu_0 M_S H_z} \]  

(2.16)

From the fact that this energy barrier has to be overcome by the thermal energy, an expression for the field \( H_z \) at which nucleation occurs can be found.

When the external field is increased further after nucleation, two scenarios are possible. Nucleation sites can be formed everywhere in the structure, eventually resulting in completely reversed magnetization. It can also be that before a field is reached that allows for this nucleation reversal, the created domain walls are moved. This is the regime our samples fall within.

2.2 Out-of-plane materials and in-plane fields

2.2.1 Relevant results from literature

Precession torque driven domain wall motion in literature

To our knowledge, precessional domain wall motion in PMA materials driven by in-plane field pulses is a topic not investigated before, and similar experiments could not be found in literature. Very relevant however, is a recent paper investigating precessional domain wall motion in in-plane materials [25]. The idea is identical to what is discussed in section 2.1.2; a magnetic field is applied perpendicular to the direction of the spins in the domain wall, resulting in a precession torque which moves the domain wall. In the paper it is successfully shown that domain walls in in-plane materials can be moved this
way. The material parameters in PMA nanowires are different from the ones in in-plane nanowires, especially the damping parameter is significantly larger in our Pt/Co/Pt systems than in the samples used in this article. Also the method to insert a domain wall in the wire is different; when an in-plane nanowire is grown in a U-shape and a field is applied along this U, a domain wall is created in the curved part. A similar trick cannot be performed on PMA nanowires, we are able to inject a domain wall by locally lowering the anisotropy by gallium ion irradiation, which will be discussed in detail in chapter 3. Still the results from this article are promising, also for the out-of-plane case.

Influence of in-plane fields in literature

Though in literature it has not been tried to move domain walls in PMA materials using in-plane fields, the influence of static in-plane fields on magnetization reversal processes has been studied. In an article by Ryu et al., reference [26], it is investigated how the velocity of a domain wall driven by current is affected by static in-plane fields. The effect turns out to be rather large; fields along the wire can even invert the direction of the motion. The presence of Dzyaloshinskii-Moriya interaction (DMI) could explain these observations, a phenomenon discussed in the next subsection.

In another paper, reference [27], the influence of in-plane fields on the domain wall velocity in field driven motion is described. A nucleation point is found, and the expansion of the domain around this point is monitored. When in-plane fields are applied, an asymmetric expansion is observed; the domain wall can be significantly slowed down or accelerated by fields in the direction perpendicular to the wall (which in a nanowire would correspond to fields along the wire). Again, DMI is necessary to explain these results.

In-plane fields do not only influence the domain wall velocity, they may also determine where an inverted domain nucleates. This was observed by Pizzini et al. [28]; in the magnetic microstructure they investigate, a domain wall nucleates at the edge of the structure, but on which edge can be controlled by an in-plane field. This can again be explained using DMI.

We have to keep in mind that in our experiments similar effects can occur besides the precessional motion. In the next subsection DMI, which apparently causes a wide range of effects in combination with in-plane fields, will be introduced.
2.2.2 Dzyaloshinskii-Moriya interaction

Introduction to DMI

The Dzyaloshinskii-Moriya interaction (DMI) was already predicted in 1958 [29, 30], but in recent years it has gained a lot of attention. One reason for this attention is the recent observation in 2009 [31–33] of a new kind of spin structure called a skyrmion. Skyrmions are vortex like structures with a size of only a few nanometer. They are stable and can be moved by currents densities in the order of $10^6$ A/m, 6 orders of magnitude smaller than required for the motion of domain walls [34]. These are all properties that are convenient in the data storage industry. Skyrmions are only stable in systems that exhibit sizable DMI.

DMI is an antisymmetric exchange interaction; the “normal” exchange interaction tries to make neighbouring spins align while DMI tries to put them at an angle with respect to each other, which is described by equation 2.17.

\[ E_{DMI} = -\mathbf{D}_{12} \cdot (\mathbf{S}_1 \times \mathbf{S}_2) \]  

(2.17)

\( \mathbf{S}_1 \) and \( \mathbf{S}_2 \) are two neighbouring spins and \( \mathbf{D}_{12} \) the DMI vector. Asymmetry is an essential ingredient for this interaction. This is schematically shown in figure 2.8, using a 3 site indirect exchange mechanism, as was done by Fert [35]. The figure shows two neighbouring spins and a third atom which interacts with these spins as well. In part (a) of this figure, two situations are shown. After spin inversion and rotation of the first situation, the second situation is obtained. These are two operations that can always be performed without changing the energy of the system, so the two situations are identical in energy. However, in part (b) the symmetry is broken by the third atom, and it is no longer possible to obtain the second situation by rotation and inversion of the first. This means that a potentially an energy difference between the two situations exists, corresponding to the DMI. This shows that asymmetry is essential for DMI, which is quantified in the fact that \( \mathbf{D}_{12} \) in equation 2.17 is proportional to \( r \times x \) [36].

In ultrathin films, DMI arises at the interface between the magnetic layer and a layer of strong spin-orbit coupling material, like platinum. A magnetic layer has of course two interfaces, the top and the bottom. If these interfaces are identical, the effects from the top and bottom exactly cancel each other. So to get a net effect from the DMI, structural inversion asymmetry is necessary. In this project Pt/Co/Pt structures are studied, which are symmetric on first sight. However, it is reported in literature
2.2 Out-of-plane materials and in-plane fields

Figure 2.8: (a) Two configurations that can be transformed into each other using only rotation and spin-inversion of the system. (b) For asymmetric systems this is not possible [37].

that effects from DMI are measured in similar samples [27]. It is argued that Co grows different on top of Pt than Pt on top of Co, so the two interfaces are not identical [38].

When the net DMI is larger than the critical value, realistically in the order of $D = 6 \text{ mJ/m}^2$, a ferromagnetic state is no longer stable and skyrmions will spontaneously form. For slightly lower values of $D$ a metastable state exists; this is regime is interesting for applications because individual skyrmions could be written (though this writing process is another technological challenge) to represent data. In our sample and in the articles cited in the previous subsection, the DMI is not strong enough to stabilize skyrmions. However, it can still affect experiment involving domain wall motion and static in-plane fields, and how this works will be explained next.

Effects of DMI

DMI prefers rotation of the magnetization in a specific way. For a system where the magnetization does not vary along the $z$ or $y$ axis, the interface DMI density in the continuous form is [39]:

$$E_{DMI} = D \left( m_z \frac{\partial m_x}{\partial x} - m_x \frac{\partial m_z}{\partial x} \right)$$  \hspace{1cm} (2.18)

$\vec{m}$ is a unit vector in the direction of the magnetization, and $m_x$ and $m_z$ are its $x$ and $z$ components, respectively. The DMI vector is replaced by a constant $z$ component, $D$. This is allowed when the sample is isotropic in the plane and the symmetry breaking
Chapter 2. Theory

is only at the $z$ surface. Assume that $D$ is positive and then examine situation (a) in figure 2.9. In the left part of the domain wall $m_z$ is negative and $\frac{\partial m_z}{\partial x}$ is positive, so the first term in equation 2.18 lowers the energy. On the right part of the domain wall both change sign, $m_z$ is positive and $\frac{\partial m_z}{\partial x}$ is negative, combined still resulting in a term that lowers the energy. In the second term $m_x$ is positive and $\frac{\partial m_x}{\partial x}$ is positive as well, so this term also results in a lowering of the energy. In the case of a Bloch wall $m_x$ and its derivative are 0, so there is no energy contribution from the DMI. This shows that DMI favours Neél walls over Bloch walls. However, when situation (b) in figure 2.9 is examined, it can be shown, using similar arguments, that the energy of this configuration is higher than in the case of a Bloch wall. So DMI favours Neél walls, but only with a certain chirality, which depends on the sign of $D$.

\begin{equation}
E/t = \pi \left( \sigma_0 \mp 2\lambda \mu_0 M_S H_x \right) - \mu_0 M_S \pi R^2
\end{equation}

The first terms represent the energy cost of the domain wall and the second terms the Zeeman energy of the reversed domain. For the domain that nucleates in the interior the in-plane field $H_x$ does not appear in the equation. This is because any change in Zeeman energy for a spin in the domain wall will be compensated by a spin on the other

This property can be used to explain the observations mentioned in the previous subsection. The orientation of the domain wall determines whether the spin Hall effect can exert a torque and in what direction, explaining the effect observed in current-driven motion [26]. High in-plane fields along the wire can force a Neél wall with the opposite chirality, thereby inverting the direction of domain wall motion.

The preferred chiral Neél walls can also explain the dependence of the nucleation position on the field using an energy argument. We assume that DMI is present in the system, so that domain walls will take a Neél configuration with a fixed chirality. Nucleation of a domain can take place either in the interior of a structure or at the edge, see figure 2.10. For both cases the energy can be written down. For the nucleation in the interior of the domain, equation 2.15 still gives the energy correctly, though the value of $\sigma_0$ is different in the presence of DMI. In the case of nucleation at the edge, equation 2.19 gives the energy for a reversed domain.

\begin{equation}
E/t = \pi \left( \sigma_0 \mp 2\lambda \mu_0 M_S H_x \right) - \mu_0 M_S \pi R^2
\end{equation}

Figure 2.9: Two Neél walls with opposite chirality. From equation 2.18 it can be deduced that configuration (a) is favoured by the DMI energy while configuration (b) is not.
side of the domain. At the edge however, the Zeeman energy from the in-plane field does not cancel out. This can cause an increase or decrease in energy, depending on the direction of the in-plane field and on which edge you are studying. Based on this model the nucleation field at the edge should linearly depend on the applied in-plane field, which is exactly what is observed. This linear behaviour is not expected to continue at sufficiently strong fields; then the assumption that we have Néel walls with a fixed chirality will become incorrect. The in-plane field for which this change in behaviour occurs would be a measure for the strength of the DMI.

![Figure 2.10](image)

**Figure 2.10:** Top view of a PMA structure with an inverted domain nucleating (a) in the interior of the structure (b) at the edge of the structure.

In the paper describing asymmetric expansion of an inverted domain under in-plane fields [27], domain wall motion is described by the creep law, which was given in equation 2.6. For the description of the asymmetric domain expansion, the creep law is adjusted; \( \alpha \), normally just a scaling constant, is assumed to depend on the domain wall energy, which in its turn depends on the applied in-plane field \( H_x \) and the strength of the DMI. Their experimental results are represented by the points in figure 2.11, which shows the the \( z \) field required to reach a certain domain wall velocity as a function of the applied in-plane field. The adjusted creep law including a DMI field is represented by the solid curve seems to described the measured points nicely.

![Figure 2.11](image)

**Figure 2.11:** Two-dimensional equi-speed contour map of \( V \) as a function os \( H_x \) and \( H_z \), as found in literature [27]. Points represent measured data, lines represent fits using the adapted creep law 2.6. The dotted line indicates the symmetric axis, corresponding to the strength of the DMI field.
2.3 Summary

This concludes the theory chapter. The systems that will be investigated were introduced in an abstract way: domain walls in nanowires with a perpendicular magnetic anisotropy. The possible types of domain walls are discussed as well. The idea to move these domain walls using in-plane fields is explained intuitively, and the collective coordinates model and simulations, that will be used to investigate this theoretically, are introduced. Last, recent results from literature about the influence of in-plane fields on domain wall motion in PMA materials are discussed. The next chapter will provide the practical background necessary to follow the results presented later; a detailed description of our samples and their production process is given. Also the measurement setup is introduced.
Chapter 3

Experimental setup

This project comprises both measurements and sample fabrication, so for both parts an experimental setup will be given. First the techniques used for samples fabrication will be introduced as well as a description of the samples that are produced. Next an introduction to the measurement equipment will follow, together with a description of the measurement routine.

3.1 Sample fabrication

3.1.1 Electron Beam Lithography

In the theory section, a model for the dynamics of domain walls in magnetic nanostrips was derived, and our next goal is to investigate such nanostrips experimentally. To manufacture these structures on a micrometer scale, a technique called electron beam lithography (EBL) is used. In this technique, a sample is first coated with a layer of resist. Irradiation by an electron beam chemically alters this resist, so it can be used to write a pattern. After this patterning, the sample is submerged in a developer, in which only the parts of the resist that have been irradiated (or only the parts that have not been irradiated, depending on the resist) dissolve. The sample is then covered with the desired material, for instance using sputtering, a technique that will be discussed in the next subsection. The sample is then cleaned with acetone in a ultrasonic bath, which removes the remainder of the resist together with the material deposited on top. What is left is the written pattern, consisting of the desired material. A schematic overview of this process is given in figure 3.1.

As a resist, polymethyl methacrylate (PMMA), dissolved in anisole is used. This is a positive (it is removed by the developer only when irradiated), high resolution resist.
Chapter 3. Experimental setup

Figure 3.1: Process of creating a structure using EBL.

To improve the lift-off process, and hence the sharpness of the edges, two layers with a different hardness are used to create an overhang structure, see again figure 3.1. The first (soft) layer consists of 495k A6 PMMA where the numbers refer to the molecular weight and the mass percentage of PMMA in the anisole respectively. The second (hard) layer consists of 950k A2 PMMA. The PMMA is spincoated on the sample, using 5000 rotations per minute for 50 seconds. After this, the sample is baked on a 150 °C hotplate for one minute to evaporate the anisole. This procedure is identical for the two different layers. The electron beam irradiation is performed using a Raith-Elphy Quantum system installed on a FEI dual beam, which has a maximum resolution of 50 nm. For the exposure a dose of 250 \( \mu \text{C/cm}^2 \) is used. Development is done using a 1:3 MIBK isopropanol mixture for 45 seconds, and with this step the template for our nanostrips is completed.

3.1.2 Sputtering

After a template for nanowires is produced using EBL, the actual magnetic material has to be deposited. To do this, a technique called sputtering is used. The sample is brought into a vacuum chamber, in our setup the base pressure is typically between \( 10^{-8} \) and \( 10^{-7} \) mbar. In this chamber also a target of the material we want to deposit is present. Argon gas is let into the chamber and a pressure in the order of \( 10^{-2} \) is reached. A large voltage (typically between 250 V and 500 V) is applied between the target and an anode ring, ionizing the argon gas and accelerating the argon ions towards the target. When the energetic argon ions hit the target, target atoms are released. When these atoms reach the sample, they condense on it, forming a thin layer. This way ultrathin layers can be grown with Ångstrom precision. A schematic overview of this process is shown in figure 3.2.

3.1.3 Gallium Ion Irradiation

A necessity for all measurements that will follow in this thesis, is being able to nucleate and pin domain walls. One method to achieve this, is gallium ion irradiation [40]. The perpendicular anisotropy in ultra thin cobalt layers is induced by the platinum/cobalt
3.1 Sample fabrication

interfaces, like described in section 2.1.1. The gallium ions damage these interfaces and make them more rough, leading to a reduction of the perpendicular anisotropy. Experimentally this is done with a focused ion beam, which is part of the FEI dual beam system. Ga\(^+\) ions with an energy of 30 keV are used and the irradiation dose is in the order of 1.0 \(\mu\)C/cm\(^2\).

When out-of-plane fields are applied to a PMA sample, the area with the lowest anisotropy will switch first. This means that, with certain field strengths, it is possible to switch only the gallium irradiated part of a sample, so a domain wall can be nucleated [41]. Another consequence of gallium irradiation is a transition between high and low anisotropy regions in the sample. The energy of a domain wall, which was given in equation 2.5, scales with \(\sqrt{AK}\), so a step in the anisotropy causes a step in the energy landscape, see figure 3.3 (a). When an external field in the \(z\) direction is applied, the Zeeman energy decreases when the domain with the preferred orientation expands. This gives a linear contribution to the energy landscape of the domain wall. If the linear contribution is small, a local minimum is formed at the anisotropy barrier and the domain wall is pinned, which is depicted in situation (b) of figure 3.3. When the applied field increases, the linear contribution increases and at a certain point the minimum disappears so the domain wall depins, see figure 3.3 (c). This approach makes it possible to perform the depinning experiments that will be described in section 3.2.2.

3.1.4 Our samples

Our structures are grown on top of silicon substrates with a top layer of 100 nm silicon dioxide, SiO\(_2\). Because we will conduct experiments involving electrical currents, it this electrically isolating top layer is necessary to prevent additional conduction paths through the substrate. The small thickness is chosen to ensure good thermal contact.
Chapter 3. Experimental setup

Figure 3.3: (a) Energy landscape for a domain wall at an anisotropy barrier. (b) When an external field is applied, the energy landscape tilts, and the domain wall is pinned. (c) When the external field is increased, the tilting increases and the domain wall depins.

between the structure and the Si substrate. This is again in view of the experiments with currents, so the generated heat can be transported away efficiently, so higher currents can be reached without burning the sample. The structures we produce have thicknesses in the order of nanometers, so small contaminations on the substrate can destroy the properties of the sample. To prevent this, the substrates are cleaned before any structuring or growth. Cleaning is done using ammonia, acetone and isopropanol, and blowing the sample dry using a N$_2$ gun.

The first step of the fabrication process is the production of “large” (20 µm wide at the smallest parts) cross structures on a substrate. In the experiment current will be send through these structures, which will generate magnetic fields according to Ampere’s law. The pattern is written using EBL and the structure consist of 14 nm titanium and 140 nm gold, deposited using e-beam evaporation. This process is similar to sputtering, but the target atoms are not released by energetic ions but by heating by an electron beam. Gold is chosen because it is a commonly used material for high frequency experiments and because it grows smooth enough to grow magnetic layers on top (this last property is proven in appendix B). The titanium acts as an adhesive layer between the gold and the SiO$_2$. A picture of one cross structure is shown in figure 3.4 (a), the large contact pads come together in the centre in a 20 µm wide structure. This central part is enlarged in figure 3.4 (b), where also the positions of the nanostrips, that will be deposited later, are shown.

The structure is covered with 10 nm SiO$_2$ deposited by radio frequent sputtering. This is similar to the process described in section 3.1.2, but now the voltage between the target and anode is alternating. This prevents a build-up of charge on the target, making the method suitable for the deposition of insulating materials. The SiO$_2$ layer forms an insulting barrier between the nanowires and the golden cross structure; this ensures that any effects on the domain wall motion in the nanowires we observe are not current induced, like the spin Hall effect or the Rashba effect. The contact pads of the golden structure should not be covered with SiO$_2$, therefore they are covered with pen before
the RF sputter deposition. After the deposition, the pen is removed using acetone, together with the SiO$_2$ grown on top.

Figure 3.4: (a) Golden cross structure on the samples, picture is taken in a probe station. (b) Zoom-in of the central part of the cross structure where the nanowires are positioned.

Next, the magnetic nanowires are fabricated on top of the cross structures. The pattern is again written using EBL, which results in wires with a width of 1 $\mu$m and a length of 10 $\mu$m. The wires consist of several layers, all deposited using sputtering. Before the layers are deposited, the samples are cleaned using an oxygen plasma to remove residual hydrocarbons, left behind by the chemical cleaning process. The first layer is 5.0 nm thick, consisting of tantalum. This material has the property that it grows very smooth, so it improves the quality of the layers and their interfaces that are grown on top. After this, 4.0 nm platinum, 0.6 nm cobalt and again 4.0 nm platinum is grown. The cobalt is the actual magnetic material, and the interfaces with the platinum favour an out-of-plane magnetization, which is needed for our experiment. To make it possible to nucleate and pin domain walls in these wires, the anisotropy of the middle part of the wires is lowered by irradiation with gallium ions. A schematic view of the complete sample is shown in figure 3.5 together with a picture of a set of nanowires. Last, to incorporate the sample in the measuring setup it is glued onto a chip carrier, which is connected to the contact pads of the sample using wire bonding.

3.2 Measurements

3.2.1 the Kerr microscope

All measurements performed in this thesis are done using a Kerr microscope, a powerful tool to investigate magnetic nanostructures. Not only can small structures be studied like in a regular microscope, also magnetism gives rise to additional contrast. This makes it for instance possible to see a domain wall move through a magnetic structure. Figure 3.5 (b) is an example of an image taken by Kerr microscope; it shows magnified nanowires,
Figure 3.5: (a) Schematic representation of the samples used in this project. (b) Typical Kerr image of a set of nanowires on top of a larger golden structure. The brighter areas in the middle of the wires are irradiated with gallium ions.

and the centres of the wires have a different brightness because the magnetization has the opposite direction there. In the software governing our Kerr microscope (a system from Evico) it is possible to select a region of interest in the microscopic image, and measure the intensity in this region. If the magnetization reverses, this can be observed clearly by a jump in the intensity.

The Kerr microscope is based on the magneto optical Kerr effect (MOKE). Linear polarized light can be seen as a superposition of right-handed and left-handed circular polarized light. These two component will have different reflection coefficients depending on the magnetization. This results in both a rotation and ellipticity change when the light is reflected from a magnetic surface. By sending light through a polarizer before it hits the sample, and through a perpendicular analyzer before it is collected, the collected light is a measure for the rotation, and hence for the magnetization. Magnetization in every direction can give rise to a Kerr effect; in the plane of the investigate sample and along the plane of incidence of the light (longitudinal MOKE), in the plane of the sample and perpendicular to the plane of incidence (transverse MOKE) and perpendicular to the plane of the sample (polar MOKE). Because we study PMA samples, we are only interested in polar MOKE. To be sensitive to this component, we ensure that the angle of incidence is close to zero. Detailed information about MOKE can be found in various sources, such as [42].

3.2.2 Depinning field measurements

In this thesis, hysteresis loops are obtained by measuring the intensity of the image of the non-irradiated parts of the nanostrips. In figure 3.6 a typical expected loop is shown. A loop starts at a field of 15 mT in the negative z direction, which saturates the magnetization everywhere in the nanostrap, situation (a). The z field is now gradually
3.2 Measurements

swept to positive values, and after every change of 0.2 mT the intensity is measured. When positive \( z \) fields are reached, the magnetization in the Ga\(^+\) irradiated area will switch at a certain moment, introducing two domain walls in the nanowire, which are pinned at the anisotropy boundaries, situation (b). When a magnetic field is increased further, the domain wall will depin at a certain moment and propagate through the wire, situation (c), which is detected by the Kerr microscope as a change in intensity. The field at which this happens is called the \textit{depinning field}. To complete the hysteresis loop, this procedure is repeated while sweeping the \( z \) field from positive to negative values. The hysteresis loops presented in this thesis are all processed. The raw data usually contains a linear slope due to the Faraday effect, which is subtracted. Also the measured intensity cannot be converted to the absolute magnetic moment, so the data is normalized in such a way that a value of -1 corresponds to magnetization in the -\( z \) direction and +1 to magnetization in the +\( z \) direction. Please note that the measurement routine described here concerns field driven domain wall motion in the conventional way, the precession torque due to in-plane fields is not yet involved.

![Figure 3.6](image)

**Figure 3.6:** Expected hysteresis loop when measuring the top part of a nanowire. Schematic representation of the magnetization in the nanowire is given for various places in the loop. The depinning field can be identified as the sharp jump in the Kerr signal.

What we are interested in is how this depinning field changes when in-plane field pulses are applied. We expect intuitively that if pulses are applied that exert a precession torque that should move the domain walls in the direction of the non-irradiated area, depinning should become easier, and therefore the depinning field should become lower. To test this experimentally, current pulses are applied through the larger golden cross structure during the measurement, which generates Oersted fields experienced by the nanowires. For the generation of the current pulses several devices are available. The
“picosecond” pulse generator can generate pulses with a rise time of 60 ps. Unfortunately our setup is not designed for high frequency experiments, and we are forced to use a “kepco” as a pulse generator with a rise time of 10 µs. Our setup is also equipped with electromagnets which can apply a static magnetic field up to 38 mT during the measurement, in any desired direction.
Chapter 4

Results from modeling and simulations

4.1 Can domain walls be moved by the precession torque?

In the theory section, the idea of precession torque driven domain wall motion is introduced and intuitively explained, but no solid evidence for this idea is provided. In this chapter, two more formal tools, the collective coordinates model and OOMMF simulations that have been introduced in section 2.1.3 and 2.1.4, are used to investigate whether precession torque driven domain wall motion is possible.

An in-plane magnetic field pulse is applied to an out-of-plane magnetized nanowire in which a domain wall is present, and the displacement of this domain wall during and after the pulse is studied. A schematic representation of the system, including the polarity of

![Figure 4.1](image_url)

**Figure 4.1:** (a) Schematic representation of the investigated system, including the definition of the domain wall displacement $q$, and the in-plane angle $\phi$. (b) Domain wall displacement during a field pulse according to OOMMF simulations (solid circles) and 1D model calculations (lines). Influence of the strength of the applied field pulse is investigated.
the domains, the direction of the spins in the domain wall, the direction of the applied field and the definition of the domain wall displacement $q$, is given in figure 4.1 (a). The nanostrips are simulated in OOMMF as rectangles with a length of 2000 nm, a width of 150 nm and a thickness of 1 nm, using material parameters matching the properties of our samples (which can be found in literature [17], supplementary information). This is a system in which Bloch walls are stable when DM interactions are not taken into account. A Bloch wall is artificially created and the system is relaxed to obtain the initial configuration. A field pulse is then applied along the length of the wire (the $x$ direction) with a duration of 1 ns and a rise and fall time of 100 ps. The influence of the strength of the applied field is investigated, it is varied from 10 to 50 mT, which are realistic values that we expect to obtain in our experimental systems. A calculation on the same system is done using the 1D model. The applied fields strengths, damping constant and gyromagnetic ratio are identical to the ones used in the simulation. The other input parameters needed are the domain wall width, $\lambda$, and the demagnetization field, $H_D$. These values can be estimated from the formula 2.4 and the in-plane angle of the domain wall according to the simulation, respectively. Note that rise and fall time are not included in these calculations. Qualitatively, there is good agreement between domain wall motion as predicted by the simulation and as predicted by the 1D calculation. Quantitatively discrepancies in the displacement of the domain wall up to 10 percent are observed. By not fixing the values for $\lambda$ and $H_D$ in the 1D calculations, but determine them by fitting with the simulations, good agreement between simulation and calculation can be achieved. Results for calculations using these parameters are shown in figure 4.1 (b) together with the simulation results. The points in the figure are simulation results, the lines the 1D model calculations.

Both simulations and 1D model show domain wall displacement, which is promising because this suggests that domain wall motion driven by the precession torque is indeed possible! The domain wall velocity is high, in the order of 100 m/s, and turns out to increase for higher applied fields. Because the precession torque is proportional to the applied field this is an expected result. Further it is observed that the position of the domain wall does not remain constant when the external field pulse ends after one nanosecond. Simulations and calculations have been performed for times up to a microsecond after the pulse (not shown here), and it turns out that the domain wall eventually moves back to its original position. Why this happens and how it can be prevented will be discussed in subsection 4.4. Experimentally this could cause additional difficulties, so is preferable to show the working principle of precession torque driven motion in an experiment that is not hindered by this behaviour.

Also, we observe that during the field pulse the domain wall velocity gradually decreases. We expect that this is because of the damping, and this is investigated further with more
4.1 Can domain walls be moved by the precession torque?

Calculations and simulations. The strength of the applied field is kept at a constant value of 40 mT this time, while the damping parameter is varied between 0.1 and 0.4. These are realistic values for the out-of-plane material [43]. Also the duration of the pulse is increased from 1.0 ns to 4.0 ns to investigate how the decrease in velocity continues. Results are shown in figure 4.2 (a), again the circles are obtained by simulation and the lines by 1D model calculations. Initially, the domain walls all move with the same velocity, irrespective of the strength of the damping, see the parts of the graphs in figure 4.2 (a) up to 0.5 ns. Then the domain wall velocity decreases, and for stronger damping this decrease is faster. Eventually, the domain wall motion comes to a halt. The physical interpretation is the following: the damping term in the LLG equation, which is proportional to the Gilbert damping parameter, makes the spins align with the applied magnetic field. When the spins in the domain wall align with the field, the precession torque decreases, and the domain wall motion comes to a halt, schematically represented in figure 4.2 (b).

In summary, both OOMMF simulations and calculations using a 1D model show domain wall motion in PMA materials by in-plane field pulses, convincingly showing the validity of this new idea. The results of the two methods are in excellent agreement, both with respect to the qualitative behaviour of the domain wall, and the quantitative values of the displacement and velocity. By varying parameters it was demonstrated that large and fast domain wall displacement can be obtained using high magnetic fields and low damping materials. Now the basics of precessional motion are understood, we can investigate if it can be extended to every type of domain wall.
4.2 Bloch walls, Néel walls and chirality

For both the 1D model calculations and OOMMF simulations in the previous section, it is assumed that the initial configuration of the domain wall is of the Bloch type. Though this is what could be argued using arguments involving demagnetization and geometry of the nanowire, as explained in section 2.1.5, we cannot be certain that this is indeed the case for our samples. DMI as introduced in section 2.2.2 for instance, could force the domain wall into a Néel configuration, and it is not straightforward to determine the orientation of the spins in the wall in our setup. Therefore, it is interesting to investigate whether the precession torque moves a Néel wall as well, and if the behaviour is similar to the Bloch wall.

\[
\begin{align*}
\dot{q} & = \frac{1}{1 + \alpha^2} \left( \frac{p}{\alpha} \right) \left( \frac{\pi \gamma}{2} (H_x \cos \phi_{\text{Bloch}} + H_y \sin \phi_{\text{Bloch}}) - \frac{\gamma}{2} H_{DBloch} \sin 2\phi_{\text{Bloch}} \right) \\
\dot{\phi}_{\text{Bloch}} & = 1 + \alpha^2 \\
\end{align*}
\]  
(4.1)

The equations of motion derived with the 1D model for a Bloch wall, again shown in equation 4.1, are adjusted to describe a Néel wall. All symbols have the same meaning as in section 2.1.3 where this equation was introduced, and the subscripts \(\text{Bloch}\) and \(\text{Néel}\) indicate whether a Bloch or Neel wall is described. The demagnetization field \(H_D\) determines whether a Bloch or Néel wall is energetically more favourable, as explained in section 2.1.5, so this parameter has to change sign. The in-plane angle can be defined in such a way that it is zero in the Néel configuration, making a comparison more straightforward.

\[
\begin{align*}
H_{DNéel} & = -H_{DBloch} \\
\phi_{Néel} & = \phi_{\text{Bloch}} - \pi/2 \\
\end{align*}
\]  
(4.2)

The rules in equation 4.2 mathematically describe these adjustments.

\[
\begin{align*}
\dot{q} & = -\frac{1}{1 + \alpha^2} \left( \frac{p}{\alpha} \right) \left( \frac{\pi \gamma}{2} (H_x \sin \phi_{\text{Néel}} - H_y \cos \phi_{\text{Néel}}) - \frac{\gamma}{2} H_{DNéel} \sin 2\phi_{\text{Néel}} \right) \\
\end{align*}
\]  
(4.3)

After rewriting the equations of motion, equation 4.3 is obtained for Néel walls, which has exactly the same form as equation 4.1. Note that to move Bloch walls fields have to be applied in the \(x\) direction while for Néel walls fields have to be applied in the \(y\) direction.
4.2 Bloch walls, Néel walls and chirality

direction. So according to the 1D model both Bloch walls and Néel walls can be moved using the precession torque, and both should behave in the same way. This is tested using OOMMF; a Néel wall is artificially inserted, the system is relaxed and an in-plane magnetic field pulse is applied perpendicular to the wire. Simulation parameters are identical to the ones in section 4.1, only the width of the nanowire is reduced to 40 nm in order to stabilize the Néel wall. The obtained results are indeed similar to the ones in figure 4.2 (not shown here), confirming that both Bloch and Néel walls can be moved using the precession torque. A small difference in velocity, in the order of 10 m/s, between the Bloch wall and Néel wall is observed, which can be explained by a different magnitude of the demagnetization field. Note that these insights provide us with an easy experimental procedure to distinguish a Bloch wall from a Néel wall: one can simply determine the direction of the in-plane field that is required to move the domain wall!

Even when it is certain that you are dealing with, for instance, a Bloch wall, four different configurations are possible by changing the polarity, $p$ of the domains and direction of the spins in the wall $c$, see figure 2.6. These four situations can be split into two groups; the left-handed and right-handed chirality, also shown in figure 2.6. By intuitively determining the direction of movement by the precession torque, as explained in section 2.1.2, it can be deduced that the domain wall should move in opposite direction for the two chiralities. In the 1D model the polarity can be reversed by changing the value of $p$ from 1 to -1, and the direction of the spins in the domain wall can be changed by taking $\phi = \pi$ as an initial condition instead of $\phi = 0$. Studying the equation of motion, it can be derived that this indeed gives an extra minus sign for the displacement. For the OOMMF simulations this can be used to test whether the domain wall displacement we observe is really due to the precession torque or caused by some artefact. Simulations starting from all four initial configuration are performed and the results are shown in figure 4.3 (a). We indeed observe that the domain wall displacement is in opposite direction for opposite chiralities, which confirms our interpretations of the results in section 4.1.

So, not only can Bloch and Néel walls be distinguished using precession torque driven motion, also chirality of the domain wall can easily be determined from the direction of domain wall motion, see figure 4.3 (b). This is advantageous over other methods to study the type of domain wall, for instance resistance measurements, because of this information about the chirality and because individual domain walls can be investigated. In studies like [26], where it is argued that Néel walls with a fixed chirality are necessary to explain their results, domain wall characterisation using the precession torque could be a very useful tool.
4.3 Adding pinning sites

As mentioned in section 4.1, the domain wall returns to its original position when the in-plane field pulse ends. For applications and experiments permanent domain wall displacement would be preferable, and also over unlimited distances. Fortunately, solutions are possible and will be introduced in section 4.4. However, these solutions would require further complication of our setup, so it is decided to show the working principle of the precessional motion in another way; using depinning experiments as described in section 3.2.2. In these experiments, the domain wall is moved by a $z$ field, but pinned at the anisotropy boundary. In-plane fields exerting a precession torque can assist the depinning. However, when the in-plane field pulse ends, the domain wall is already at a remote distance of the pinning site driven by the $z$ field, and cannot return to its original position. This approach is different from moving a domain wall purely with an in-plane field pulse, as is modeled in the previous subsections. Therefore, the simulations and calculations done up to now cannot give quantitative expectations for the experiments, and an modification of the model is necessary.

To mimic the depinning experiment, an artificial pinning site is added to the one dimensional model, for details see appendix A.3. The pinning site is modelled as an finite harmonic potential well, as in reference [44], where the maximum pinning strength is calculated from the anisotropy [40]. A schematic representation of the energy landscape for the domain wall is shown in figure 4.4 (a). In the calculation, the domain wall starts at the position of the pinning site, and a $z$ field is slowly increased until the wall gets depinned. This $z$ field will give a linear contribution to the energy landscape, see figure 4.4 (b). It is now investigated how the value of the $z$ field at which this happens is altered when in-plane field pulses are applied.
4.3 Adding pinning sites

The depinning field is indeed decreased by pulses applied in the correct direction, schematically shown in figure 4.4. By “correct direction” is meant a pulse that exerts a precession torque which tries to move the domain wall in the same direction as the $z$ field. It turns out that pulses in the opposite direction decrease the depinning field as well, though this effect is significantly smaller than the effect from the correctly oriented pulses. This can intuitively be explained in the following way; the domain wall is moved over a small distance in the wrong direction by the pulse. After this, the $z$ field will move the domain wall back, and it will hit the steepest point of the pinning site with a certain velocity. With this velocity it is easier to overcome the pinning than starting from rest, as depicted in figure 4.4 (d).

During the experiments, hysteresis loops are measured, which actually include two depinning events; sweeping the $z$ field from negative to positive and back. During the experiment, we will fix the direction of the spins in the domain wall in a certain direction using a static in-plane field, so this is done in the calculations as well. Note that the chirality of the domain wall will be opposite when the $z$ field is swept up or down. For instance, sweeping up the pulses will be in the correct direction to help the depinning, while sweeping down the pulses are in the wrong direction. From this, two expectations about the experimentally obtained loops can be made; first the pulses will decrease the coercivity because both going up and down the depinning field will be decreased, see figure 4.5 (b). Also, because the change in depinning field is equal in size going up and going down, we expect a shift of the loops, which is shown in figure 4.5 (a). In this figure it can be seen that the shift is indeed 0 mT when no pulses are applied and that
the loop shift to the positive or negative side depending on the direction of the applied pulse. Also we observe that the direction of this shift is inverted when when the initial in-plane angle is changed from 0 to $\pi$, corresponding to a change in chirality.

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure4.5.png}
\caption{Predicted effects on (a) the shift of hysteresis loops and (b) the coercive fields from pulses in the depinning experiment by the one dimensional model.}
\end{figure}

4.4 Preventing the domain wall from returning to its original position

Though the model and simulations in section 4.1 provide us with a deep understanding of precession torque driven domain wall motion, also a disadvantage is discovered: when the external magnetic field pulse ends, the domain wall will return to its original position. The spins are affected by the total field, which consists not only of the external field pulse, but also has contributions from the anisotropy and the demagnetization. So when the external field is switched off, there is still an effective field present, which is under a certain angle with the spins. This effective field will exert a precession torque on spins, making the domain wall move back. For future experiments on this topic, and certainly for possible future applications this moving back of domain walls would be a bottleneck, so in this subsection two solutions are suggested.

4.4.1 Pinning

Pinning sites can prevent a domain wall from moving when a field or current tries to drive it, so this could also be used to prevent a domain wall from moving back. Simulations have shown that this indeed works for in-plane materials [25]. However, experimentally this could be problematic, because the pinning will have to meet specific requirements. The pinning has to prevent the domain wall from moving back, but not from moving in the first place. This sets conditions for the pulse shape of the applied field, the strength of the applied field and the geometry of the sample and the pinning strength. Another
point to consider is the separation between the different pinning sites; this should be less than the distance a domain wall can travel during one pulse. The resolution reached with gallium irradiation, the method we use to create pinning sites in PMA material, is in the order of 10 nm [40]. Whether this is small enough will depend on the material parameters, like the damping constant. If this distance is too large, one would have to rely on natural pinning sites, which has the disadvantage that they cannot be controlled.

4.4.2 Using fields in the other in-plane direction

Pinning might be a good solution, but it cannot be guaranteed that it will work experimentally, so now another solution is introduced. This solution is based on the precession torque. Suppose a Bloch wall is moved by a field pulse in the $x$ direction. As was shown by the model and simulations in section 4.1, the velocity slowly decreases and the movement eventually comes to a halt. This is because of the damping which makes the spins in the domain wall align with the applied field, and after a while a Néel wall is created on which the $x$ field does not exert a precession torque. This Néel wall in its turn could be moved by a field in the $y$ direction. So if we do not go back to zero field after applying the $x$ field pulse, but immediately apply a $y$ field pulse, the domain wall motion could be continued in the desired direction. The spins will align with the applied field, which will result in a Bloch wall. If the field is turned off now, the spins in the domain wall are already aligned with the effective field, so there is no precession torque to drive the domain wall back!

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{figure4_6.png}
\caption{Application of a field pulse in the $y$ direction directly after the pulse in the $x$ direction can prevent the domain wall from returning to its original position.}
\end{figure}

This idea has been tested by the model and by simulations, for which the results are shown in figure 4.7. When the $x$ field is applied the domain is moved, just as before, and the motion comes to a halt after a few nanoseconds. In figure 4.7 (a) the graph for the in-plane angle $\phi$ is included, which shows that the domain wall indeed transforms
from a Bloch wall to a Néel wall during this pulse. Now a $y$ field is added. Because it is experimentally not feasible to switch the $y$ field on at the exact moment when the $x$ field is switched off with sub-nanosecond precision, a period in which the $x$ and $y$ field are applied simultaneously is included. The total field is now again at an angle with the spins in the domain wall, and exerts a torque, so the wall is moved during this period as well. This is followed by a period with only a $y$ field in which, as expected, the domain wall is again moved, until a Bloch wall is formed. In the last part of the simulation and calculation the applied fields are switched off. It is observed that, contrary to the results with only an $x$ field pulse like in figure 4.2, the position of the domain wall remains constant. Even if the calculations are extended to times in the order of microseconds (not shown here), the domain wall does not move back to its original position, proving that our solution is effective.

![Figure 4.7: Solution for moving back of the domain wall; an additional in-plane field pulse is applied. (a) results obtained using the collective coordinates model, (b) results obtained using OOMMF. The time and displacement are different from each other and from results obtained in section 4.1. This is because different material parameters are used, which is merely due to the chronological progress of the project and we expect to observe similar behaviour when the parameters from section 4.1 are used.](image)

A variation on this solution would be the use of a rotating field. Figure 4.8 shows a simulation of a domain wall driven by a rotating field with an amplitude of 40 mT and a period of 400 ns for one period. This way the distance that can be travelled by the domain wall is unlimited; it continues as long as the rotating field is applied. The fundamental limit for the domain wall speed is set by how fast the magnetization can follow the applied field, which will be in the order of nanoseconds. In practice the speed will be limited by the maximum frequency of the rotating field. An advantage is that using a rotating field may be easier in practice than two individual pulses applied directly after each other.

An observant reader may have noticed the following problem; if there are two domain walls present with equal chiralities, an $x$ field drives them in the same direction, but when the spins are aligned with this field and Néel walls are formed, they will have opposing
Preventing the domain wall from returning to its original position

Figure 4.8: Domain wall motion as a function of time when a rotating in-plane field is applied. The graph is obtained using OOMMF simulations.

chiralities. The subsequent $y$ field will therefore drive them in opposite directions. For samples with strong DMI this can be solved. Without applied external fields, the two domain walls will experience an opposing effective field. Using this fact, it is possible to design a sequence of pulses that moves the domain walls in a synchronous way. Figure 4.9 gives an example of such a sequence; after this sequence the left domain wall is back at its original position, while the right domain wall has moved to the right. A next sequence of $+H_y$, $+H_x$ and $-H_y$ pulses would now move the left domain wall to the right. Experimentally this is beyond the scope of this thesis, so no verification using the model or the simulations is performed. However, if research on precessional motion would be continued, this would be an interesting approach.

Initial configuration:

After applying: $H$

After applying: $-H$

After relaxation:

Figure 4.9: Intuitive scheme for irreversible precession torque driven motion for multiple domain walls. The division of the nanowire into squares is merely a guide to the eye.
4.5 Summary

Both the collective coordinates model and OOMMF simulations show that a domain wall in a PMA nanostrip can be moved by the precession torque from an in-plane field pulse. The influence of the strength of the field pulses and Gilbert damping parameter is investigated and understood. It was demonstrated that this approach is suitable for both Bloch and Néel type domain walls, and can even be used to characterize the type of domain wall. An undesired feature of this method for domain wall motion is the tendency of the domain wall to return to its original position after the pulse ends. Solutions for this problem have been suggested and checked using the model. To mimic the depinning experiments that are performed in this thesis, a pinning site is added to the model and calculations of the depinning field are performed, giving both qualitative and quantitative predictions for the experiments.
Chapter 5

Depinning field experiment

After the modeling of precession torque driven domain wall motion in the previous chapter, our experimental results are now presented. In the experiment we investigate the change in depinning field when pulsed or static in-plane fields are applied, and discuss whether this change is caused by the precession torque.

5.1 Effects of current pulses

5.1.1 Reference experiment

To study the effect of in-plane fields, we first need a reference measurement without in-plane fields. Figure 5.1 (a) shows a typical set of nanowires that is studied during the experiments, and the “top” and “bottom” part are indicated for future reference. In this figure, two regions of interest (ROIs) are indicated by the red and blue rectangles, these are the areas in which the intensities of the image are measured during the measurements. A step in this intensity corresponds to magnetization reversal and consequently to domain wall depinning. These steps are clearly visible, in figure 5.1 (a) for instance, the top parts of the wires are brighter than the bottom parts, corresponding to a different direction of the magnetization. The setup for the measurements is described in section 3.2.2.

Typical hysteresis loops are shown in figure 5.1 (b). The large steps in the graphs, where the Kerr signal switches between -1 and +1, indicate the switching of the top (or bottom) part of the nanowire and therefore the depinning of the top (or bottom) domain wall. In this report, two features of obtained hysteresis loops are studied; the first is the shift of the loops, which is the position of the middle of the loop along the horizontal axis. In this reference experiment we expect this shift to be zero, because
the experiment is symmetric for the depinning of an up-to-down domain wall and an
down-to-up domain wall. If a shift would be measured, this would indicate an offset in
the applied $z$ field. The average of the two switching fields in the left part of the graph
is $-9.0$ mT and the average of the switches in the right part is $9.3$ mT. The average
shift is therefore $0.15$ mT. During the measurement, the $z$ field is increased with steps
of $0.2$ mT, so the measured shift is not significant, as expected. What also needs to be
added to this minimal uncertainty of $0.2$ mT, are thermal fluctuations. The difference
in depinning field during the up sweep and down sweep is determined for 12 different
domain walls (not shown here), and this difference amounts to $(-0.2 \pm 0.4)$ mT. The
difference is not significant, as expected, and the standard deviation gives a indication of
the magnitude of the thermal fluctuations. This explains, for instance, why the switches
of the two loops in figure 5.1 (b) overlap on the left part of the graph, but not on the
right part. Note that an overlap of these two loops is not forced by symmetry arguments;
the loops concern domain walls pinned at different positions, and local pinning sites can
significantly affect the measured coercive field!

![Figure 5.1: (a) Typical Kerr image of the nanowires investigated in this thesis. The
red squares indicate the regions of interest that are measured to determine the switching
fields of both the top and bottom domain walls. (b) Typical hysteresis loops of both
the top and bottom part of a nanowire.](image)

The second feature of the loops we are interested in is the coercive field, which amounts
to $9.2$ mT in this reference experiment. The official definition of coercivity is “the
intensity of the magnetic field needed to reduce the magnetization of a ferromagnetic
material to zero after it has reached saturation”. Because we are performing a depinning
experiment, this is not the physical interpretation of the switching fields in our loops.
So when the coercive field is mentioned in this thesis, this refers to the average value of
the two switching fields in one hysteresis loop.
Now these reference hysteresis loops are measured and analyzed, it can be investigated how they change when in-plane magnetic fields are applied.

5.1.2 Results with pulses from Kepco

Current pulses are applied to the sample in the configuration as depicted in figure 5.2, generating Oersted fields.

![Figure 5.2: Measurement configuration used in this section.](image)

This turns out to strongly affect the depinning, the depinning field changes in the order of 10 mT for the largest applied pulses. Figure 5.3 (a) shows hysteresis loops with and without applied current pulses. We observe that the effect has the opposite sign for the two domain walls in one wire, which would make sense if they had the same chirality. However a reason for this is not evident. Dzyaloshinskii-Moriya interactions can favour a certain chirality for Néel walls, as discussed in section 2.2.2. But domain walls should be the Bloch type to be affected by a precession torque from Oersted fields in the $x$ direction that are used here. The hysteresis loops have more unexpected features; most evident, a switch can occur at the “wrong” side of zero on the horizontal axis, so unfavourable with respect to the Zeeman energy. Even if the short pulses assist, or even cause, depinning, the Zeeman energy is the driving force to propagate the domain wall through the rest of the wire. Another feature that stands out in the loops is the small “extra switch” in figure 5.3 (b), indicated by the blue, dashed circle. The wires consist of only one magnetic layer, which should result in a hysteresis loop with only one step. Also this feature is not observed when the pulses are applied in the opposite direction, though a similar result would be expected on symmetry grounds. An explanation can be found in the Kerr images taken during the switching process; for strong current pulses the switching turns out not to be abrupt. During the measurement the $z$ field is increased by steps of 0.2 mT. With every step the domain wall moves over a distance in the order of a micrometer. Apparently during the measurement of the loop with the “extra switch” the domain wall paused somewhere within the region of interest.

Looking at the Kerr images something else is noticed; the domain wall does not come from the Ga$^+$ irradiated area, but from one of the far ends of the wire! To show this,
we present images from movies taken with the Kerr microscope. First examine figure 5.4; the images in this figure are taken without applying current pulses. The sample is saturated, the field is brought back to 0 mT, an image for background substraction is taken, and the measurement is started. At the start of the measurement, the nanowire can only be distinguished vaguely because of the background substraction. Normally the external \( z \) field would be swept gradually from -15 mT to 15 mT, but the origin of the the inverted domain cannot be observed this way. Therefore the \( z \) field is pulsed (the pulse duration is 3 ms) so a domain wall cannot propagate through the whole wire during one pulse. In the three right images it can be seen that the inverted domains originate in the middle of the wire, the Ga\(^+\) irradiated area, the dashed blue circles indicate where this can be noted.

A similar measurement while applying 1.0 A current pulses gives very different results, see figure 5.5. The left image in this figure is again taken at the start of the measurement. Note despite the clear contrast, the magnetization is saturated! This contrast is the result of taking an image for background substraction at 0 mT applied field, and as can be seen from the loops in figure 5.3 one half of the nanowire has already switched in this situation. In the right image it can be seen that the inverted domains nucleate at the top end of the nanowires, again indicated by the blue dashed ellips.
5.1 Effects of current pulses

Taking all this into account, we conclude that the observed effects cannot be explained by precession torque assisted depinning. In the following sections we try to explain the aforementioned effects and propose a new measuring configuration to observe effects from the precession torque.

5.1.3 Out-of-plane field component

A possible explanation for the observed difference in behaviour between two domain walls in one wire could be an out-of-plane component of the generated magnetic field. An infinite plane with a constant current density only results in an in-plane Oersted field. Our cross structure through which the current is send, has a width of 20 µm. Though this is three orders of magnitude larger than its distance from the nanowire (about 10 nm) deviations from the infinite plane approximation might occur. A \( z \) component of the Oersted field could explain all the observed unexpected features discussed in the previous section. There should be a difference between hysteresis loops measured on the left and right part of the wire, for the \( z \) component of the Oersted field will have an opposite sign. If the \( z \) component of the Oersted field is large enough, it could also explain the switching at the “wrong side of zero”, because the total \( z \) field can now have another direction than the \( z \) field applied with the external magnet.

The Ga\(^+\) irradiated area is exactly above the middle of the golden strip, where any \( z \) component of the Oersted field should disappear on symmetry grounds. So, potentially large, contributions in the \( z \) direction from the Oersted field, can arise at the ends of the nanowires, but no at the centre, which could explain domain walls originating from the edges. Also note the horizontal wires in figure 5.5. These wires were not studied up to now, because in a first attempt we assume the domain walls to have a Bloch configuration, and in the horizontal wires the generated in-plane field would not exert a precession torque on Bloch walls. However, the fact that the top and bottom wires have a different contrast in this image indicate that one wire has switched at 0 mT while the
other one has not, which would be perfectly compliant with the out-of-plane Oersted contribution hypothesis.

So a \( z \) component of the Oersted field could explain all unexpected observations. However, before doing the experiments, we intuitively expected that this component would be negligible in our samples, because of their close resemblance to an infinite plane. To reliably determine the correctness of the \( z \) component hypothesis, we perform a numerical calculation of the \( z \) component of the field, as a function of the position along the cross structure, 10 nm above it, see figure 5.6 (a). The nanowire is 10 \( \mu m \) long, so it would experience the field that are calculated from -5 to 5 \( \mu m \) along the \( x \) axis. In figure 5.6 (a) it can be seen that in this range the \( z \) component of the Oersted field varies from -10.7 to 10.7 mT. This a huge contribution, the difference between the highest and lowest value is 68% of the \( x \) field component. This is large enough account for our observations.

![Figure 5.6: Numerical calculations of the generated Oersted field of a 1.0 A current through a 150 nm thick and (a) 20 \( \mu m \) wide (b) 50 \( \mu m \) wide golden wire, together with a cartoon of the situation.](image)

5.1.4 Wider cross structures

If we are indeed dealing with a \( z \) component of the Oersted field, it should be possible to reduce the effect by changing the geometry of the sample. The calculations in the previous section are repeated for a cross structure with a width of 50 \( \mu m \) instead of 20 \( \mu m \), the results are shown in figure 5.6 (b). The general behaviour is the same as for the 20 \( \mu m \) wide structure; there is a significant \( z \) component which increases when approaching the edge of the gold. However, over a distance of 10 \( \mu m \) along the \( x \) axis the change in \( z \) component is only 27% of the field in the \( x \) direction. This is an improvement compared to the 20 \( \mu m \) wide structure, where the variation in \( z \) component is 68% of the \( x \) field. Note that is an improvement of a factor 2/5, which is expected, for we changed
5.1 Effects of current pulses

the width with the same factor. This means that for experiments on samples with wider cross structures, a smaller contribution from $z$ components of the Oersted field would be expected.

We continue producing samples with 50 µm wide cross structures. The difference in shift of the hysteresis loop on both sides of the wire is plotted in figure 5.7. For comparison, the measured points at 1.0 A for the 20 µm wire are shown as well. The shift is clearly smaller for the 50 µm wide wire, 1.5 mT at maximum. A part of this difference is due to the reduced current density when the same total current flows through a wire with a larger cross section. This can be seen in figure 5.7; the field strength is decreased for the wider structure with the same total current, by a factor of $2/5$ to be exact. The remainder of the difference is because the wider structure is more resemblant to the case of an infinite plane, because the distance between the nanowires and the edges of the goldens structure is larger. A kink is observed in the graph for the wider wire, only for currents of 1.0 A and larger the difference in shift between the left and right domain wall start to occur. An explanation would be that this kink occurs when the $z$ component at the end of the nanowire is so large that a domain wall no longer depins at the Ga$^+$ irradiated area, but nucleates at the edge. This hypothesis cannot be proven from our data; the switches are faster than the time resolution of the camera on the Kerr microscope, and the origin of the inverted domain cannot be observed. An option to investigate this, would be using a pulsed $z$ field, like in section 5.1.3.

![Figure 5.7](image.png)

**Figure 5.7:** The difference in shift of the hysteresis loop between the two dw’s in one wire as a function of the static magnetic field along the wire.

The fact that changing the width of the cross structure changes the shift of the hysteresis loops in the way that is predicted by the numerical calculations, confirms that a $z$
component of the Oersted field is the correct explanation for our observations. Though increasing the width of the cross structure could even further decrease this $z$ field component, this is not an ideal solution. The ratio between the $x$ component and $z$ component will increase, but the absolute value of the $x$ field will be lower as well. This $x$ field exerts the precession torque and should be as large as possible, so high currents would be necessary to compensate for its decrease. An alternative solution will be discussed in the next subsection.

5.1.5 A solution: measurements on the middle wire

In the previous subsection, we measured a large effect from current pulses on the depinning field which cannot be explained by the precession torque, but can by a $z$ component of the generated Oersted field. This hypothesis is supported by numerical calculations and an additional experiment with a different sample geometry. The Oersted $z$ field contribution makes it more difficult to conclusively observe effects from the precession torque, so it would be preferred if the measurement setup could be adapted to exclude this effect.

![Figure 5.8: Measurement configuration used in this section.](image)

In figure 5.6 it can be seen that the $z$ component of the Oersted field vanishes at $x = 0 \ \mu m$, so above the middle of the cross structure. If the nanowire is positioned along the current direction instead of perpendicular to it, it is possible to place a complete wire in an $z$ component free region (or at least without a difference between the top and bottom part of the wire), see the setup in figure 5.8. Measurements are performed using this new geometry. Again the hysteresis loops are affected by the applied pulses, see figure 5.9. When pulses are applied the coercive field decreases by several mT. Also the loops shift when pulses are applied and this shift is opposite for the top and bottom part of the nanowire. We hope to observe that the precession torque from the generated in-plane field assist the depinning. According to the calculations in section 4.3 this would indeed
lead to a decrease in coercivity and a shift of the hysteresis loops. So are the features in
figure 5.9 the effect from the precession torque that we are looking for? To answer this
question, we need to know the orientation of the spins inside the domain walls. Are we
dealing with Bloch walls or Néel walls? And what is their chirality? Answers to these
questions are necessary to predict which observations are expected from the precession
torque.

It is decided to apply static in-plane fields using an external magnet to control the ori-
entation of the spins in the domain walls. This way the type of domain wall and its
chirality are known. So applying pulses in the sample geometry introduced in this sub-
section, together with a static in-plane field, will enable us to make a direct comparison
between our experimental results and calculation. This will lead to a substantiated con-
clusion about the presence of precession torque effects in our setup. However, in order to
exclude new effects caused by the static fields, first some reference experiments without
current pulses will be discussed in the next section.

5.2.1 Static fields along the nanowire

To compare the results from experiment and model, the configuration of the spins inside
the domain wall has to be known, as discussed in the previous subsection. The precession
torque will only have an effect if the spins in the domain walls are perpendicular to the
pulsed field, so a static field along the nanowire is used to ensure we have this desired
configuration. In this section results are presented that are obtained using only these
static in-plane fields, but no pulses, see figure 5.11 (a).

Initially this experiment is performed mainly for completeness, but we observe an in-
teresting phenomenon. The hysteresis loops are shifted, for example, during the upward
sweep of the z field the depinning field is increased, while during the downward sweep

\[ \text{Figure 5.9: Hysteresis loops for both DWs on the nanowire with applied pulses of:} \]
\[ \text{(a) -1.0 A (b) 0.0 A (c) 1.0 A.} \]
the depinning field is decreased, see figure 5.11 (b). This shift is in the opposite direction for the top and bottom domain wall in the same wire. Reversing the direction of the external in-plane field reverses the direction of the shifts, compare figure 5.11 (b) and (c).

In figure 5.12 the difference in shift between the top and bottom domain wall is plotted as a function of the strength of the static in-plane field. The difference in shift turns out to be proportional to the in-plane field, and this trend reproduces for samples grown in different sessions (in figure 5.12 the results for 3 different samples are shown). This observation is not expected from precession torque theory; those effects should occur when the system is not in equilibrium, so a pulsed field would be necessary instead of a static field. Also one of the arguments used in subsection 5.1.2 applies again; a field along the nanowire only gives rise to a precession torque if the walls are of the Bloch type. A different behaviour between the two walls in one wire can only be explained if their chirality is the same, and there is no interaction that would force this on Bloch walls. Also the effect cannot be explained by a misalignment of the applied field. Misalignment may indeed give rise to an out-of-plane field contribution in the order of a mT might exist, but it is not possible that this out-of-plane component differs in the order of a mT for the two domain walls, that are only a µm apart.
5.2 Effects of static fields

Figure 5.12: The difference in shift of the hysteresis loop between the two dw’s in one wire as a function of the static magnetic field along the wire. Wires on three samples grown in a different session are investigated.

A possible explanation can be found by comparing our results to a similar experiment in literature [28]. They claim that an in-plane field can change the energy of a domain wall, for systems with DMI. This makes it favourable for an inverted domain to nucleate at a certain side of the magnetic structure, controlled by the direction of the in-plane field, like discussed in section 2.2.2. We study domain walls that depin from a Ga\textsuperscript{+} irradiated area while they study domains that nucleate at the edge of the structure. Instead of calculating the energy difference between one or no domain wall, we should calculate the energy difference between a domain wall in the irradiated or in the non-irradiated area. A similar model could be applied to our situation, and explain our observations if we assume that the irradiation reduces the DMI. This is a reasonable assumption because the DMI in our systems is an interface effect, and according to literature [45] ion irradiation increases the interface roughness of Co/Pt.

We now consider a simplified system to demonstrate how a change in DMI together with an in-plane field can affect the depinning field. Normally the pinning of a domain wall at an anisotropy barrier is described by an energy barrier formed by the difference in domain wall energy because of the different anisotropy. Now a different DMI strength in combination with an in-plane field can give an additional contribution to this energy barrier, the situation is schematically shown in figure 5.13. We now show that the contribution from the DMI together with the in-plane field is significant compared to the original energy barrier. It is assumed that the DMI in the Ga\textsuperscript{+} irradiated area has disappeared completely and the anisotropy is reduced from $K = 1.5 \times 10^6\, \text{J/m}^3$ to
$K = 1.3 \times 10^6 \text{J/m}^3$, which would be reasonable according to [40]. The other parameters are kept the same and it is assumed that the DMI in the non-irradiated areas is strong enough to enforce a Néel wall with a certain chirality. The domain wall is pinned because of the change in anisotropy, which results in an energy barrier of $4\sqrt{AK_{\text{non-irradiated}}} - 4\sqrt{AK_{\text{irradiated}}}$. Using $A = 1.6 \times 10^{-11} \text{J/m}$ [17], this corresponds to $0.0014 \text{J/m}^2$ of domain wall. The difference in energy due to the in-plane field can be calculated assuming the spins in the irradiated area can align with the field, while the spins in a domain wall outside that area can not because of the DMI. The energy difference is therefore $2E_{\text{Zeeman}}$ with $E_{\text{Zeeman}} = \mu_0 \lambda M_S H_x$ per m$^2$ of domain wall, which corresponds to a difference of $0.041 \text{J/T}$. A magnetic field of 10 mT gives therefore rise to an energy difference of 30% of the original energy barrier. Because this energy difference that scales with the in-plane field only exist when the DMI forces the domain wall in the opposite direction of the field, different depinning field would be observed when sweeping the $z$ field up or down, resulting in a shift of the hysteresis loop which increases for stronger fields. These are all crude approximations, but they confirm that reduced DMI in the Ga$^+$ irradiated area is a possible explanation for the observed trend.

![Energy landscape for a domain wall in a nanowire with a Ga$^+$ irradiated area.](image)

Figure 5.13: Energy landscape for a domain wall in a nanowire with a Ga$^+$ irradiated area.

At the time this thesis was written, DMI was intensively investigated by many research groups. Our setup makes it possible to do quantitative measurements on a DMI related phenomenon in an easy way. So, even though it is not related to the precession torque, two additional experiments to further investigate this effect are carried out and discussed in the next two subsections.

### 5.2.2 Higher in-plane fields

In the article that inspired us to explain our observations with a combination of DMI and Zeeman energy [28], in-plane fields up to 250 mT are used. This is almost an order of magnitude larger than the fields we use, and for comparison it would be useful to obtain
5.2 Effects of static fields

data for this high fields as well. Our experiment is repeated with higher in-plane fields. The typical behaviour of the top and bottom domain wall in one wire is shown in figure 5.14 (a) and 5.14 (b), respectively. For fields up to 50 mT the depinning field varies linearly with the strength of the static in-plane field. For the top DW, the depinning field increases for positive x fields when sweeping the z field from negative to positive, while the depinning field decreases when sweeping the z field from positive to negative. This corresponds to a shift of the hysteresis loop. For the bottom DW, figure 5.14 (b), the change in depinning field is the other way around, meaning that the hysteresis loops shift in the opposite direction. This difference in behaviour is exactly what was observed before in figure 5.12.

This linear trend does not continue for x fields higher than 50 mT. The depinning field decreases for these high fields to about 2 mT at 300 mT, independent of the direction of the in-plane field or the sweep direction of the z field. This measurement is done for 6 different domain walls, and what is remarkable is that the maximum depinning field is not found at the same in-plane field value. This suggests that the behaviour is very sensitive, for instance, to small random pinning sites.

These results cannot be described by the simple model introduced in the previous sub-section. We still believe that DMI should be used in the description, for the behaviour of the domain walls is chiral. Probably more effects should be included in the model: there are more possible configurations for the spins in the domain wall except for aligned with the DMI field or antiparallel to it. In literature [39] the in-plane angle of the domain wall is calculated as a function of the applied field and the DMI strength, and this should be incorporated in the model. Also the possibility that at certain field strengths it becomes more favourable for a domain to nucleate at the edge of the wire than for the domain wall to depin from the irradiated area, should be taken into account.

For future modeling, another approach could be chosen as well. In a recent study by Je et al. [27], asymmetric expansion of a magnetic domain under applied in-plane fields are studied. Because this experiment involves in-plane fields, DMI and comparable samples, we believe that the effect they observe is strongly related to the observations by Pizzini and our own observations. They explain their results using an adjusted creep law where DMI is taken into account, as explained in section 2.2.2. This creep law could be fitted nicely with their experiments depicted in figure 2.11. This figure shows the z field required to move a domain wall at a certain speed as a function of the x field. Intuitively this quantity plotted as a function of the in-plane field is similar to the depinning field that we plot as a function of the in-plane field. When figure 5.14 is compared to figure 2.11 a resemblance in shape is observed, which might be an indication that we could use a similar approach to describe our results. During creep motion, the domain wall
is kind of depinning the entire time, so from that point of view it may be intuitively understandable that our depinning experiment results would follow the creep law. We believe a model developed from this starting point could therefore be valuable in the description of our observations.

5.2.3 Variation of the Ga$^+$ dose

An advantage of our system over the experiment in literature [28] is that we can vary the properties of the region the domain wall comes from, and therefore control the energy difference for the domain wall. Though our current model needs to be extend before reliable predictions can be made, as discussed in the previous subsection, we have intuitive expectations. The basic hypothesis is that Ga$^+$ irradiation reduces the DMI, which makes it easier to align with an external field. This results in a difference in Zeeman energy for a domain wall in the two different regions. We expect that a larger difference in DMI therefore results in a larger difference in Zeeman energy and so a stronger dependence on the external field. It is interesting to check this experimentally because it could show a decrease in DMI by ion irradiation and confirm our model.

A new sample is produced with nanowires which are irradiated with different Ga$^+$ doses. The difference in shift of the loops between the two domain walls is shown for different Ga$^+$ doses in figure 5.15 (a). The doses are given in units of $\mu$C/cm$^2$, and 1.0 $\mu$C/cm$^2$, for our irradiated areas with a surface of approximately 1 $\mu$m $\times$ 1 $\mu$m corresponds to $6.25 \times 10^4$ gallium ions.

A dose of 1.0 $\mu$C/cm$^2$ turned out to be the minimal value required to do a depinning measurement. For lower doses the difference in anisotropy is too small to act as a pinning site. On the other end of the spectrum we have a dose of 5.0 $\mu$C/cm$^2$, for which the irradiated area is non-magnetic; when measuring a region of interest in this area, no
hysteresis is observed. The graph for the 5.0 µC/cm² dose has a significantly higher slope than the graph for the 1.0 µC/cm² dose. The 1.5 µC/cm² and 2.5 µC/cm² dose graphs lay in between the two extremes. To show this in a more transparent way, the slopes of the graphs in figure 5.15 (a) are plotted as a function of the Ga⁺ dose in figure 5.15 (b). The slopes are determined using linear fits, though a model should confirm that linear behaviour is expected. We clearly see that the influence of external in-plane field increases with the Ga⁺ dose, which is in agreement with our intuitive expectations. This is an indication that Ga⁺ irradiation indeed decreases the DMI of our samples.

Though the graphs in figure 5.15 are noisy and have large error bars, the difference in slope when increasing the dose from 1.0 µC/cm² to 5.0 µC/cm² is significant, as can be seen from the slopes plotted in figure 5.15 (b). Therefore it would be interesting to improve this experiment; measurements at more different doses could be performed, and to reduce the noise in the graphs averages could be taken over more nanowires (in the results presented here every data point is an average over two nanowires). This would be advisable for further research on DMI and Ga⁺ irradiation.

In summary; we observe a clear chiral dependence of the depinning field on static in-plane fields. This is evidence of DMI in our seemingly symmetric samples, which is in agreement with recent observations in literature. Also we present preliminary results of the effect of varying the Ga⁺ dose. These results indicate that the irradiation reduces the DMI, as would be intuitively expected, and is an interesting topic for further research.

5.3 The complete experiment; pulses and static fields

The depinning field measurement is now repeated both with field pulses and a static field to control the spin configuration in the domain walls, see the geometry in figure 5.16. If
the orientation of the spins in the domain wall is known, as well as the direction of the pulses, quantitative predictions can be made using the 1D model, as were presented in section 4.3. Now the results that should match these calculations are presented in the next subsection.

![Figure 5.16: Measurement configuration used in this section.](image)

### 5.3.1 Final results

Measurements are performed, and are compared to the calculations in figure 4.5. Figure 5.17 (a) shows the coercive field as a function of the current pulse $I$. Our model predicted a decrease of the coercive field with current, because the current pulses assist the depinning and therefore lower the depinning field. We indeed observe this decrease in coercive field, independent of the static field. However, the shape of the graph does not perfectly match the predicted behaviour, see figure 4.5 (b), it looks rather parabolic. This parabolic behaviour could indicate another effect which reduces the coercivity: heating. The increase in temperature is proportional to the power, which is proportional to the current squared. To test whether heating plays a significant role in our setup, a new measurement is performed; the pulse strength and duration are kept constant but the number of pulses per second is varied. The results are shown in figure 5.17 (c). A clear dependance is observed; the more pulses per second the smaller the coercivity becomes. This indicates that heating significantly reduces the coercivity in our experiments (in general, a frequency of 1000 pulses per second is used), and no conclusions regarding the precession torque can be drawn based on the measurements in figure 5.17.

Because the coercivity cannot provide any prove for effects by the precession torque, a feature independent of heating should be investigated. Because heating changes the depinning field when sweeping up by the same amount as when sweeping down, it cannot account for a shift of the loops. Also we have a quantitative prediction about the shift of the loops from the 1D model, so this feature is perfect to investigate next! In figure 5.18, the shift of the hysteresis loops is plotted as a function of the strength of the applied
5.3 The complete experiment; pulses and static fields

Figure 5.17: (a) Coercive field as a function of the in-plane pulse strength, measured together with static in-plane fields. (b) The coercive field as a function of the number of pulses per second.

pulse. This is done with the static field applied in two opposite directions. Note that the static field determines the chirality of the domain walls, so reversing this field should reverse the direction of the shift, like predicted in figure 4.5 (a). However we can clearly see that this is not the case in the measurement. Unfortunately, we have to conclude, on symmetry grounds, that the observed features do not originate from the precession torque.

Figure 5.18: Shift of hysteresis loops as a function of the strength of the applied in-plane field pulses for two different static in-plane fields.

One may wonder whether the static fields we apply (up to 38 mT) are strong enough to control the orientation of the domain wall, because this seems to be in contradiction with the results and interpretation in section 5.2; there it is argued that spins in the domain wall keep aligned favourable to the DMI field up to external fields of 50 mT. If this interpretation is correct, 38 mT would indeed be too small to control a domain
Chapter 5. Depinning field experiment

wall in the non-irradiated area of the wire, but for this experiment only the orientation before the domain wall gets depinned is of importance, so when the domain wall is inside the irradiated area. Based on similar samples that were investigated in our group [17] we expect that fields of around 20 mT are enough to switch, for instance, from Bloch to Néel wall. The strength of the in-plane fields can therefore not explain the absence of precession torque effects.

The graphs in figure 5.18 clearly have a slope, so the in-plane field pulses do have an effect. However, the origin of this effect is still unclear at this moment. An explanation would again be a z component of the Oersted field. If the current is not homogeneously distributed over the cross structure, but takes certain, random, preferential paths, this z component could arise, even though care has been taken to prevent this. In the next subsection, we will not continue this discussion about what we do see in figure 5.18, but investigate why we do not see effects from the precession torque.

5.3.2 Influence of the rise time

The 1D model predicted a change in depinning field of several mT due to the precession torque caused by in-plane field pulses, but this was not observed in the experiments. We revise the model trying to explain this discrepancy. The key parameter turns out to be the rise time of the pulse. All calculation presented so far were performed assuming a rise time of 100 ps, comparable to pulses from the “picosecond pulse generator” we intended to use at the start of this project. Eventually, the measurements were performed using a “kepco” with a rise time of 10 µs, and this turns out to greatly affect the magnitude of the precession torque. Figure 5.19 shows the difference in depinning field caused by the precession torque as a function of the rise time. For rise times of 1 ns and smaller, the effect is indeed significant; in the order of several mT. For rise times of 10 ns or longer, the effect is smaller than 0.2 mT, which is the uncertainty in our setup, so this is too small for us to measure. The fact that we do not measure any precession torque effects using pulses with a rise time of 10 µs is in agreement with these new calculations. Another consequence is that the movement of a domain wall by the precession torque of a rotating field like described in section 4.4.2, only works in theory, for perfect systems without pinning sites.

Why the rise time is critical for the precession torque can be understood intuitively. The precession is torque calculated by \( \gamma \vec{M} \times \vec{H} \). The largest value is obtained when \( \vec{H} \) has its maximum value, and \( \vec{M} \) and \( \vec{H} \) are orthogonal to each other. So large torques are obtained for short rise times, when the applied field reaches it maximum value before the spins in the domain wall are aligned with the field. Alignment with the applied field
5.3 The complete experiment; pulses and static fields

Figure 5.19: (a) Calculations varying the rise time are performed, which is a property depending on the pulse generator. (b) Change in the depinning field by in-plane field pulses as a function of the rise time of the pulse. In this calculation a pulse of 10 mT is applied in the correct direction for 10 ns excluding rise and fall time.

by damping is a process which takes places on the timescale of nanoseconds, as can be seen in figure 4.2 (b). This is indeed the timescale of importance for the rise time as we observe in figure 5.19, confirming the intuitive interpretation.

Last, the 1D model is used to investigate whether there is a way to circumvent the necessity of these very short rise times. The limiting effect is that the spins align with the pulsed field, which is governed by the damping term in the LLG equation. Therefore it is plausible that varying the damping parameter could influence the rise time of the pulse that is necessary. Varying the damping parameter is experimentally possible as well, using a different magnetic material. To investigate the effect of damping, the dependence of the change in the depinning field on the rise time is plotted for a wide range of damping parameters in figure 5.20. We observe that a larger effect from the precession torque is expected when the damping is smaller, in agreement with figure 4.2. For our samples $\alpha = 0.1$, approximately, which limits the maximum change in depinning field to about 5 mT, which can easily be detected in our setup. However, for all values of the damping parameter, the change in depinning field by the precession torque becomes negligible for rise times larger than 10 ns.

Another parameter that can be varied experimentally is the strength of the DMI, which can be done by changing growth conditions or materials. DM interactions are probably of influence in our system, as they are necessary to explain the observations in section 5.2, but up to now they have not been included in the collective coordinates model or simulations. According to literature [39], DMI can be interpreted as an additional magnetic field only felt inside the domain wall, $H_{DMI}$. This is how it is incorporated in the model; as a constant additional field which is aligned with the domain wall in its
initial configuration. Calculated plots of the change in depinning field versus the rise time of the pulse for values of $H_{DMI}$ ranging from 1.5 mT (or 0 when the results of figure 5.19 are taken into account) to 1600 mT are shown in figure 5.20. This includes all realistic values, because the range includes the symmetric case of no DMI, ±250 mT in the case of Pt/Co/AlOx [26], the strength needed to stabilize skyrmions (±450 mT) [22] and the maximum strength for which the ferromagnetic state is stable (±700) mT. It is not possible to enhance the precession torque for longer rise times by adding this DMI field. In general, the change in depinning field decreases when the strength of the DMI increases. The explanation for this is similar to the explanation presented in the previous paragraph, concerning the damping parameter; the precession torque is maximum when the magnetization is perpendicular to the effective field. The DMI field is part of this effective field, and the spins are aligned with this field in the initial configuration. So for increasing DMI field, the effective field becomes more and more aligned with the spins in the domain wall, and the precession torque becomes smaller. The ideal nanowire for precession torque driven domain wall motion would therefore have a DMI field of less than 100 mT, to get a sizable effect from the precession torque. Note that eventually, a DMI field of more than 0 mT would be convenient, so eventually the idea presented in figure 4.9 can be applied.

![Figure 5.20](image)

**Figure 5.20:** Change in depinning field as a function of rise time; investigating the effect of (a) the damping constant (b) the DMI field.

In conclusion, no significant effect from the precession torque on the depinning field are observed in our experiments. To explain this result, our 1D model is extended with a new parameter; the rise time of the pulse. This parameter turns out to be of major importance; to observe effects from the precession torque, pulses with a rise time of 1 ns or less are necessary. This would require a different setup and sample design suited for high frequencies. This is beyond the scope of this project, but in the outlook chapter the possibilities will be discussed in more detail.
Chapter 6

Conclusion

Both the theoretical study of precession torque driven domain wall motion and the experimental investigation of our samples have let to new insights about the subject. We start with the conclusions from the theoretical part.

The effect of an in-plane field pulse on a domain wall in a PMA nanowire was theoretically studied using a one dimensional model and OOMMF simulations. The results of both methods are in excellent agreement with each other and describe the same behaviour of the domain wall; the field pulse exerts a precession torque on the domain wall resulting in domain wall motion. Using parameters corresponding to our samples, we estimate a domain wall velocity in the order of 100 m/s. This velocity increases linearly with the strength of the field pulse. The domain wall motion is a non-equilibrium phenomenon; as soon as the spins in the domain wall are aligned with the new effective field, the domain wall stops moving. This aligning is governed by the damping term of the LLG equation, and increasing the Gilbert damping parameter makes the domain wall motion stop sooner.

When the field pulse ends, the new total field (including effects from anisotropy and exchange) exerts a precession torque moving the domain wall back to its original position. This can be prevented in two ways. First, it is possible to add small pinning sites and make the rise time and fall time of the pulse asymmetric. This makes the precession torque at the start and end of the pulse different in strength, so it is possible for the domain wall to move past the pinning sites when the pulse is started, but not when the pulse is ended. Another possibility would be, after the movement by the pulse has stopped, directly applying an in-plane field pulse in an other direction, which again would exert a precession torque to move the domain wall further. This can be done using individual pulses or using a rotating field.
Chapter 6. Conclusion

How an in-plane pulse influences depinning from a pinning site by an out-of-plane field is also investigated using the one dimensional model, because this resembles our experiment. The pulses turn out to lower the depinning field, so when hysteresis loops are measured, the coercive field should decrease when pulses are applied. The amount with which the depinning field changes is dependent on the chirality of the domain wall. If a measurement is performed while the orientation of the spins in the domain wall is fixed by an static in-plane field, the chirality is different during the upward sweep and the downward sweep of the hysteresis loop, because the polarity is inverted. So the pulses should shift the measured hysteresis loops.

Performing experiments, the first conclusions we drew were about additional effects. The Oersted fields used as in-plane field pulses turned out to have a significant out-of-plane component as well. Care has to be taken that experiments are performed in such a geometry that this out-of-plane component is zero. When the in-plane field pulses are applied, the coercive field decreases. However, additional experiments show that heating is a plausible explanation for this effect. So to draw conclusions about the effect from the precession torque, we have to look at the shift of the hysteresis loops, because this cannot be caused by heating.

The shift of the hysteresis loops, as predicted by the one dimensional model, is not observed. This is a clear indication that we are not observing effects from the precession torque. Further investigation using the model shows that only pulses with a rise time of 1 ns or less should significantly affect the depinning field. This explains our observations, and both sample design and measuring setup should be altered to be suitable for high frequencies in order to observe precessional motion.

Using only static fields, we do observe shifts in the hysteresis loops. Though these shifts are not due to the precession torque, they can also not be explained by simple arguments as heating or misalignments. Similar effects were reported in literature for other experiments, and explanations are given using the Dzyaloshinskii-Moriya interaction. Our results can be described by the same models with adjustments. The advantage of our experiment is that the depinning field is easy to measure, and by varying the Ga⁺ dose for the irradiated area, we have an extra parameter we can tune. So refining these measurements could give a valuable contribution to the ongoing research on DMI.
Chapter 7

Outlook

As discussed in chapter 5.3.2, the key requirement to make further progress in the search for precessional domain wall motion, is being able to apply pulses with a rise time in the order of 1 ns or less. This requires a change in setup and sample design. For the samples this would mean that the golden cross structure should be replaced by a coplanar wave guide. A coplanar waveguide consists of three conducting lines next to each other, where the middle one is the guideline for the pulse, and the outer two are grounded. If the magnetic nanowires are now placed of the guideline, pulses with an ultra fast rise time can be applied to them. To prevent distortion of the pulses before they reach the sample, they should be applied to the sample using a high frequency probe, instead of a chip carrier. A schematic overview of this new high frequency setup is drawn in figure 7.1. The depinning experiments in this thesis can now be repeated, and an effect of several mT from the precession torque would be expected.

Figure 7.1: Alternative setup to apply in-plane field pulses to nanowires, suitable for high frequencies. (a) Top view of the necessary new sample design, a coplanar waveguide with the nanowire on top of the guideline. The distances $s$ and $w$ can be used to tune the impedance of the structure. (b) A high frequency probe station, including a Kerr microscope and external magnets.
The final goal of this study is to build a field driven racetrack memory, so after the working principle of the precession torque has been demonstrated by the depinning experiments, we can focus on actual domain wall motion driven only by in-plane field pulses. To do this, it would be convenient first study a domain wall with minimal pinning. To obtain such a situation, we can again nucleate a domain wall using Ga irradiation, and then using a short magnetic pulse in the $z$ direction to move the wall away from the pinning site. To make a racetrack with multiple domain walls, it would be necessary that depinning from an anisotropy boundary can be achieved by the in-plane pulses. To do this we probably need to increase the precession torque, for instance by lowering the damping by choosing other materials. Also we need a sizeable DMI to fix the chirality of the domain walls, which can be achieved by varying the growth parameters.

Our expectation is that, though we did not manage to do it in this project, it will be possible to produce a precession torque driven racetrack memory in the future, if pulses with a short rise time are used. But it is unlikely this will be used in the computer industry. The extreme fast rise time that is necessary puts strict limitations to the design of the structure and to the pulse generator to be used. We expect that this disadvantage outweighs the advantages over other mechanisms to drive domain walls. But like pointed out in the introduction, one cannot predict how technology will evolve over the years; suppose that high frequency pulses become a commonly used phenomenon, precession torque driven domain wall motion might suddenly become interesting for industry again.

In research, precessional domain wall motion is of use despite the high frequency requirements. This is because domain walls motion is still extensively studied, but because of the small dimensions of the walls it is difficult to measure the orientation of the spins in the walls directly. This is however of great interest, it determines for instance whether the walls can be moved by the spin Hall effect and it gives information about the controversial DM interactions. Using precessional motion the orientation of the domain wall follows simply from the direction of the field pulse that is necessary for the movement. This way it can be easily determined whether we have a Bloch configuration, a Néel one, or a combination of both. Even the chirality can be determined, and the measurement can be done for an individual domain wall, great advantages over for instance using a domain wall resistance measurement.
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Appendix A

Detailed derivations 1D model

In section 2.1.3 equation 2.10 and 2.11 are given without any derivation. This is because the derivation is long and does not provide much insight. However it might be useful for the reader that wants to check equation 2.11 or for someone who is performing a similar calculation. Therefore the detailed derivation is given in this appendix. In this derivation it is assumed that at $x$ is minus infinity the spins are in the minus $z$ direction ($\theta = \pi$), and at $x$ is plus infinity the spins are in the plus $z$ direction ($\theta = 0$). The spins in the centre of the domain wall point along the minus $y$ direction and the external field is applied in the plus $x$ direction. The in-plane angle is defined to be zero if the spins point in the minus $y$ direction and $\pi/2$ if the spins point in the plus $x$ direction, see figure 2.4. Though these assumptions are now made, the end result can easily be modified for other polarities $p$, chiralities, initial domain wall configurations and directions of the applied field.

A.1 The static situation

The shape of the domain wall is determined by two opposing contributions. The first one the exchange interaction, which makes neighbouring spins align. This favours the domain wall being as stretched out as possible; the wider it is the smaller the angle between two adjacent spins is. On the other hand, the anisotropy makes the spins to be along the easy axis, in this case the $z$ direction. This contribution favours small domain walls, because when the domain wall is smaller, less spins are not pointing along the $z$ axis. The energy of the system can be expressed in the following way by integrating the energy density over the volume:
Appendix A. Detailed derivations 1D model

\[ E = S \int_{-\infty}^{\infty} \left[ A \left( \frac{\partial \vec{m}}{\partial x} \right)^2 - K (\vec{m} \cdot \hat{z})^2 + K_d (\vec{m} \cdot \hat{x})^2 \right] \, dx \]  
(A.1)

In equation A.1 \( E \) is the total energy, \( S \) is the surface of the cross section of the wire, \( A \) is the exchange stiffness, \( \vec{m} \) is the direction of the magnetization, \( K \) is the easy axis anisotropy constant and \( K_d \) the hard axis anisotropy constant. It is allowed to multiply with \( S \) instead of integrating over the surface, because the system is assumed to be one dimensional. The \( K_d \) term lifts the degeneracy between the the \( x \) and \( y \) direction, if the constant is positive Bloch walls are favoured, if it is negative Néel walls are favoured. Generally the magnetization is defined in the following way, where \( M_S \) is the saturation magnetization:

\[ \vec{M} = M_S (\sin \theta \sin \phi, -\sin \theta \cos \phi, \cos \theta) \]  
(A.2)

However, in the static case we may assume that the magnetization is completely in the easy plane. We assume, as stated in the introduction that the we have a Bloch wall with \( \phi = 0 \), then:

\[ \vec{M} = M_S (0, \sin \theta, \cos \theta) \]  
(A.3)

Combining equation A.3 with equation A.1, gives:

\[ E = S \int_{-\infty}^{\infty} A \left( \frac{\partial \theta}{\partial x} \right)^2 - K \cos^2 \theta \, dx \]  
(A.4)

A differential equation for \( \theta \) can be found using the Euler-Lagrange equations.

\[ \frac{\partial}{\partial t} \left( \frac{\partial L}{\partial \dot{x}} \right) = \frac{\partial L}{\partial x} \]  
(A.5)

Applied to our situation this results in:

\[ A \frac{\partial^2 \theta}{\partial x^2} = K \cos \theta \sin \theta \]  
(A.6)

Infinitely far from the domain wall the magnetization is perfectly out-of-plane. This gives the boundary conditions:
A.1 The static situation

\[ \theta (x = -p\infty) = \pi \]
\[ \theta (x = p\infty) = 0 \]
\[ \frac{\partial \theta}{\partial x} (x = -\infty) = 0 \]
\[ \frac{\partial \theta}{\partial x} (x = \infty) = 0 \]  

(A.7)

Here \( p = \pm 1 \) and represents the polarity. If the system has the polarity as depicted in figure 2.4 its value is +1, if the polarity is reversed its value is -1. Solving equation A.6 with these boundary conditions results in an equation for the shape of the domain wall.

\[ \cos \theta = \tanh \frac{px}{\lambda} \]  

(A.8)

When it is assumed that the centre of the domain wall is not at position \( x = 0 \) but at \( x = q \), the generalized form of the equation becomes:

\[ \cos \theta = \tanh \frac{p(x - q)}{\lambda} \]  

(A.9)

Here \( \lambda = \sqrt{A/K} \) is called the domain wall width. In figure A.1 \( \theta \) is plotted as a function of the position along the wire using this solution. The shape is as expected: the magnetization is in the minus \( z \) direction in the left side of the wire, in the plus \( z \) direction in the right side of the wire and changes gradually at the position of the domain wall (in this case \( x = 0 \)).
Appendix A. Detailed derivations 1D model

A.2 Applying in-plane fields

The general method we use to derive the equations in this section was developed by A. Thiele [19] and we follow steps like in literature [44], where it was used to derive a 1D model for current driven domain wall motion. To derive equations 2.11, the LLG equation is used, see equation A.10. The terms in this equation are numbered for future reference.

\[
\frac{\partial \vec{M}}{\partial t} = -\gamma \vec{M} \times \left( \vec{H}_{\text{ext}} + \vec{H}_{\text{eff}} \right) + \frac{\alpha \vec{M}}{M_s} \times \frac{d\vec{M}}{dt} \quad (A.10)
\]

Each of these terms can be written in the form \(-\gamma \vec{M} \times \vec{H}\). So we have to find the correct form of \(H\) for each term. Term 2 already has the correct form, so

\[
\vec{H}_2 = H_x \hat{x} + H_y \hat{y} \quad (A.11)
\]

The third term has the correct form as well. This term is the sum of the exchange interaction, the out-of-plane anisotropy and the demagnetization field. This last term determines which in-plane angle is most favorable.

\[
\vec{H}_3 = 2A \frac{\partial^2 \vec{M}}{\partial x^2} + H_K \frac{M_z}{M_S} \hat{z} - H_D \frac{M_x}{M_S} \hat{x} \quad (A.12)
\]

The first term is less clear, but the reader is invited to compute \(-\gamma \vec{M} \times \vec{H}_1\) and see that the first term of the LLG equation can indeed be recovered in this way. The same holds for term 4, the damping term.

\[
\vec{H}_1 = -\frac{1}{\gamma} \frac{d\vec{M}}{dt} \quad (A.13)
\]

\[
\vec{H}_4 = -\frac{\alpha}{\gamma} \frac{d\vec{M}}{dt} \quad (A.14)
\]

Now we can define a total field which is zero:

\[
\vec{H}_{\text{tot}} = \vec{H}_1 + \vec{H}_2 + \vec{H}_3 + \vec{H}_4 = 0 \quad (A.15)
\]
From this it follows that

\[- \vec{H}_{\text{tot}} \cdot \frac{\partial \dot{M}}{\partial x} = 0 \tag{A.16} \]

\[- \vec{H}_{\text{tot}} \cdot \frac{\partial \dot{M}}{\partial \phi} = 0 \]

Analyzing the units, these values can be interpreted as a force and a torque density respectively. Because \( \vec{H}_{\text{tot}} \cdot \frac{\partial \dot{M}}{\partial x} = \sum_{i=1}^{4} \vec{H}_{i} \cdot \frac{\partial \dot{M}}{\partial x} \), we can evaluate this force term by term.

\[ f_1 = - \left( -\frac{1}{\gamma} \dot{M} \times \frac{d \dot{M}}{dt} \right) \cdot \frac{\partial \dot{M}}{\partial x} \]

\[ = 1/\gamma \left( \frac{d \dot{M}}{dt} \times \frac{\partial \dot{M}}{\partial x} \right) \cdot \dot{M} \]

\[ = -1/\gamma \left( \frac{\partial \dot{M}}{\partial \phi} \times \frac{\partial \dot{M}}{\partial \theta} \cdot \dot{M} \right) \frac{d \phi}{dt} \frac{\partial \theta}{\partial q} \tag{A.17} \]

where we have used that \( dx = -dq \).

\[ f_2 = - (H_x \ddot{x} + H_y \ddot{y}) \cdot \frac{\partial \dot{M}}{\partial x} \]

\[ = - (H_x \ddot{x} + H_y \ddot{y}) \cdot \frac{\partial \dot{M}}{\partial \theta} \frac{\partial \theta}{\partial q} \tag{A.18} \]

\[ f_3 = - \left( 2A \frac{\partial^2 \dot{M}}{\partial x^2} + H_K \frac{M_z}{M_S} \ddot{z} - H_D \frac{M_x}{M_S} \ddot{x} \right) \cdot \frac{\partial \dot{M}}{\partial x} \]

\[ = -A \frac{\partial}{\partial x} \left( \frac{\partial \dot{M}}{\partial x} \right)^2 - \left( H_K \frac{M_z}{M_S} \ddot{z} - H_D \frac{M_x}{M_S} \ddot{x} \right) \cdot \frac{\partial \dot{M}}{\partial x} \]

\[ = -A \frac{\partial}{\partial x} \left( \frac{\partial \dot{M}}{\partial x} \right)^2 + \left( H_K \frac{M_z}{M_S} \ddot{z} - H_D \frac{M_x}{M_S} \ddot{x} \right) \cdot \frac{\partial \dot{M}}{\partial \theta} \frac{\partial \theta}{\partial q} \tag{A.19} \]
Before we continue, we need to work out some of the parameters appearing in these equations. First we examine the direction of the magnetization and its derivatives.

\[
\begin{align*}
\hat{M} &= (\sin \theta \sin \phi, -\sin \theta \cos \phi, \cos \theta) \\
\frac{\partial \hat{M}}{\partial \phi} &= (\sin \theta \cos \phi, \sin \theta \sin \phi, 0) \\
\frac{\partial \hat{M}}{\partial \theta} &= (\cos \theta \sin \phi, -\cos \theta \cos \phi, -\sin \theta) \\
\left(\frac{\partial \hat{M}}{\partial \theta}\right)^2 &= \cos^2 \theta \sin^2 \phi + \cos^2 \theta \cos^2 \phi + \sin^2 \theta \\
&= \cos^2 \theta + \sin^2 \theta \\
&= 1 \\
\left(\frac{\partial \hat{M}}{\partial \theta}\right)^2 &= \sin^2 \theta \cos^2 \phi + \sin^2 \theta \sin^2 \phi \\
&= \sin^2 \theta \\
\frac{\partial \hat{M}}{\partial \theta} \cdot \frac{\partial \hat{M}}{\partial \phi} &= \sin \theta \cos \theta \sin \phi \cos \phi - \sin \theta \cos \theta \sin \phi \cos \phi + 0 \\
&= 0 \\
\left(\frac{\partial \hat{M}}{\partial \phi} \times \frac{\partial \hat{M}}{\partial \theta}\right) \cdot \hat{M} &= (- \sin^2 \theta \sin \phi, \sin^2 \theta \cos \phi, -\sin \theta \cos \theta \cos^2 \phi - \sin \theta \cos \theta \sin^2 \phi) \cdot \hat{M} \\
&= (- \sin^2 \theta \sin \phi, \sin^2 \theta \cos \phi, -\sin \theta \cos \theta \cos^2 \phi - \sin \theta \cos \theta \sin^2 \phi) \cdot (\sin \theta \sin \phi, -\sin \theta \cos \phi, \cos \theta) \\
&= - \sin^3 \theta \sin^2 \phi - \sin^3 \theta \cos^2 \phi - \sin \theta \cos^2 \phi \\
&= - \sin^3 \theta - \sin \theta \cos^2 \theta \\
&= - \sin \theta
\end{align*}
\]

(A.21)

We also need to compute \(\frac{\partial \theta}{\partial q}\). For this, the shape of the domain wall is needed, which was derived for the static case.
\[ \cos \theta = \tanh \left( \frac{x - q}{\lambda} \right) \tag{A.22} \]

\[
\begin{align*}
\frac{\partial \theta}{\partial q} &= \frac{\partial \theta}{\partial \cos \theta} \frac{\partial \cos \theta}{\partial q} = \left( \frac{\partial \cos \theta}{\partial \theta} \right)^{-1} \frac{\partial \cos \theta}{\partial q} \\
&= -\left( \frac{1}{\sin \theta} \right) \frac{\partial \cos \theta}{\partial q} \quad \tag{A.23}
\end{align*}
\]

For simplicity, we now introduce a new variable \( z = \frac{x-q}{\lambda} \).

\[
\frac{\partial \cos \theta}{\partial q} = \frac{\partial \cos \theta}{\partial z} \frac{\partial z}{\partial q} = -\frac{1}{\lambda} \frac{\partial \tanh z}{\partial q} = -\frac{1}{\lambda} \left( 1 - \tanh^2 z \right) = \frac{1}{\lambda} \left( \cos^2 \theta - 1 \right) = -\frac{1}{\lambda} \sin^2 \theta \quad \tag{A.24}
\]

So we find:

\[
\frac{\partial \theta}{\partial q} = -\frac{1}{\sin \theta} \cdot -\frac{1}{\lambda} \sin^2 \theta = \frac{1}{\lambda} \sin \theta \quad \tag{A.25}
\]

Note that for the opposite polarity we would have started with an extra minus sign in equation A.22 and therefore have ended up with a minus sign in the final expression.

Now all “force densities” can be integrated over the whole wire.

\[
F_1 = \int dV f_1 = -\frac{1}{\gamma} S \int_{-\infty}^{\infty} dx \left( -\sin \theta \right) \frac{\partial \theta}{\partial q} = -\frac{1}{\gamma} S \int_{\theta(x=\infty)}^{\theta(x=-\infty)} d\theta \left( \sin \theta \right) = 2S \int_{-\infty}^{\infty} dt \tag{A.26}
\]

Here we used that \( dx = -dq \) and we used boundary conditions that follow from the polarity we assumed. If the polarity is reversed, the integration boundaries switch, leading to an additional minus sign.
\[ F_2 = \int dV f_2 \]
\[ = -S \int_{-\infty}^{\infty} dx \left( H_x \dot{x} + H_y \dot{y} \right) \cdot (\cos \theta \sin \phi, -\cos \theta \cos \phi, 0) \]
\[ = S \int_{\theta(x=\infty)}^{\theta(x=-\infty)} d\theta H_x \cos \theta \sin \phi - H_y \cos \theta \cos \phi \]
\[ = 0 \quad (A.27) \]

\[ F_3 = \int dV f_3 \]
\[ = S \int_{-\infty}^{\infty} dx \left[ -A \frac{\partial}{\partial x} \left( \frac{\partial \hat{M}}{\partial x} \right)^2 + \left( \frac{H_K M_z}{M_S} z - \frac{H_D M_x}{M_S} \right) \cdot \frac{\partial \hat{M}}{\partial \theta} \frac{\partial }{\partial q} \right] \quad (A.28) \]

This integral can be divided into three smaller integrals.

\[ F_{3,1} = S \int_{-\infty}^{\infty} dx \left[ -A \frac{\partial}{\partial x} \left( \frac{\partial \hat{M}}{\partial x} \right)^2 \right] \]
\[ = -SA \left( \left. \left( \frac{\partial \hat{M}}{\partial x} \right)^2 \right|_{x=\infty} - \left. \left( \frac{\partial \hat{M}}{\partial x} \right)^2 \right|_{x=-\infty} \right) \quad (A.29) \]
\[ = 0 \]

Here it is used that the magnetization is constant at infinity, so its derivative is zero.

\[ F_{3,2} = S \int_{-\infty}^{\infty} dx H_K M_z \hat{z} \cdot \frac{\partial \hat{M}}{\partial \theta} \frac{\partial }{\partial q} \]
\[ = SH_K \int_{-\infty}^{\infty} dx \cos \theta \cdot -\sin \theta \frac{\partial }{\partial q} \]
\[ = SH_K \int_{\theta(x=\infty)}^{\theta(x=-\infty)} d\theta \cos \theta \sin \theta \]
\[ = 0 \quad (A.30) \]

Here it is used that \( \frac{M_z}{M_S} = \cos \theta \), as can be seen from equation A.21, and \( dx = -dq \).
\[ F_{3,3} = S \int_{-\infty}^{\infty} dx - H_D \frac{M_x}{M_S} \hat{x} \cdot \frac{\partial \hat{M}}{\partial \theta} \frac{\partial \theta}{\partial q} \]
\[ = -SH_D \int_{-\infty}^{\infty} dx \sin \theta \sin \phi \cdot \cos \theta \sin \phi \frac{\partial \theta}{\partial q} \]
\[ = SH_D \sin^2 \phi \int_{\theta(x=\infty)}^{\theta(x=\infty)} d\theta \cos \theta \sin \theta \]
\[ = 0 \] (A.31)

\[ F_3 \text{ is therefore equal to zero as well.} \]

\[ F_4 = \int dV f_4 \]
\[ = S \int_{-\infty}^{\infty} dx - \alpha/\gamma \left( \frac{\partial \hat{M}}{\partial \theta} \right)^2 \left( \frac{\partial ^2 \hat{M}}{\partial \theta \partial q} \right)^2 \frac{dq}{dt} - \alpha/\gamma \hat{M} \frac{\partial \hat{M}}{\partial \phi} \frac{\partial \theta}{\partial q} \frac{d\phi}{dt} \]
\[ = S \int_{-\infty}^{\infty} dx - \alpha/\gamma \cdot 1 \cdot \frac{1}{\lambda^2} \sin^2 \theta \frac{dq}{dt} \]
\[ = \frac{\alpha S dq}{\gamma \lambda} \int_{\theta(x=-\infty)}^{\theta(x=\infty)} d\theta \sin \theta \]
\[ = -2 \frac{\alpha S dq}{\gamma \lambda} \] (A.32)

Here it followed from equation A.21 that the second term is zero, and we used \( \frac{\partial \theta}{\partial q} = \frac{1}{\lambda} \sin \theta \Rightarrow dx = -dq = -\frac{\lambda}{\sin \theta} d\theta \) (with a minus sign for different polarity). In this derivation the values of \( \theta \) at infinity are used again, so if the polarity is changed this would result in an extra minus sign as well, so the expression is unchanged for different polarity. Now, as stated before, the sum of all these forces should be equal to zero.

\[ \sum F = 0 \]
\[ -2 \frac{\alpha S dq}{\gamma \lambda} \frac{dq}{dt} \pm \frac{2S d\phi}{\gamma} \frac{d\phi}{dt} = 0 \]
\[ \frac{\alpha dq}{dt} - p\lambda \frac{d\phi}{dt} = 0 \] (A.33)

This is the first of the two equations stated in 2.11. The \( p \) stands for polarity, and is 1 for the polarity we used in the derivation and -1 for the opposite polarity. Note that in this derivation no assumptions about the chirality or configuration of the domain wall are made.
To find the second equation mentioned in 2.11 we are now going to look at the “torque densities” exerted on the domain wall by the various field terms.

\[
\tau_1 = -\left(-\frac{1}{\gamma} \dot{M} \times \frac{d\dot{M}}{dt}\right) \cdot \frac{\partial \dot{M}}{\partial \phi}
\]

\[
= \frac{1}{\gamma} \frac{d\dot{M}}{dt} \times \frac{\partial \dot{M}}{\partial \phi} \cdot \dot{M}
\]

\[
= \frac{1}{\gamma} \frac{\partial \dot{M}}{\partial \theta} \times \frac{\partial \dot{M}}{\partial \phi} \cdot \dot{M} \frac{\partial \theta}{dq} \frac{dq}{dt}
\]

(A.34)

\[
\tau_2 = -(H_x \dot{x} + H_y \dot{y}) \cdot \frac{\partial \dot{M}}{\partial \phi}
\]

(A.35)

\[
\tau_3 = -\left(2A \frac{\partial^2 \dot{M}}{\partial x^2} + H_K \frac{M_z}{M_S} \dot{z} - H_D \frac{M_x}{M_S} \dot{x}\right) \cdot \frac{\partial \dot{M}}{\partial \phi}
\]

(A.36)

\[
\tau_4 = \frac{\alpha}{\gamma} \frac{d\dot{M}}{dt} \cdot \frac{\partial \dot{M}}{\partial \phi}
\]

\[
= \frac{\alpha}{\gamma} \left(\frac{\partial \dot{M}}{\partial \phi} \frac{d\phi}{dt} \cdot \frac{\partial \dot{M}}{\partial \phi} + \frac{\partial \dot{M}}{\partial \theta} \frac{d\theta}{dt} \cdot \frac{\partial \dot{M}}{\partial \phi}\right)
\]

(A.37)

The next step is again to integrate these densities over the whole wire.

\[
T_1 = \int dV \tau_1
\]

\[
= \frac{1}{\gamma} S \frac{dq}{dt} \int_{-\infty}^{\infty} dx \left(\sin \theta \right) \frac{\partial \theta}{dq}
\]

\[
= \frac{1}{\gamma} S \frac{dq}{dt} \left[\theta(x=\infty) - \theta(x=-\infty)\right] \frac{\partial \theta}{dq} - \sin \theta
\]

(A.38)

\[
= \frac{2pS dq}{\gamma} \frac{dq}{dt}
\]
A.2 Applying in-plane fields

\[ T_2 = \int dV \tau_2 = -S \left( \int_{-\infty}^{\infty} dx H_x \hat{x} \cdot \frac{\partial \hat{M}}{\partial \phi} + \int_{-\infty}^{\infty} dx H_y \hat{y} \cdot \frac{\partial \hat{M}}{\partial \phi} \right) \]

\[ = -S \left( H_x \int_{-\infty}^{\infty} dx \sin \theta \cos \phi + H_y \int_{-\infty}^{\infty} dx \sin \theta \sin \phi \right) \]

\[ = -S \left( H_x \int_{\theta(x=\infty)}^{\theta(x=-\infty)} d\theta - p \lambda \cos \phi + H_y \int_{\theta(x=-\infty)}^{\theta(x=\infty)} d\theta - p \lambda \sin \phi \right) \]

\[ = -S \lambda \pi (H_x \cos \phi + H_y \sin \phi) \] (A.39)

Here we have used that \( dx = -\frac{\lambda}{\sin \theta} d\theta \), which can be seen from equation A.25.

\[ T_3 = \int dV \tau_3 = -S \left( \int_{-\infty}^{\infty} dx 2A \frac{\partial^2 \hat{M}}{\partial x^2} \cdot \frac{\partial \hat{M}}{\partial \phi} + \int_{-\infty}^{\infty} dx H_K \frac{M_S}{M_S} \hat{z} \cdot \frac{\partial \hat{M}}{\partial \phi} + \int_{-\infty}^{\infty} dx -H_D \frac{M_S}{M_S} \hat{x} \cdot \frac{\partial \hat{M}}{\partial \phi} \right) \]

\[ = 0 + 0 + S H_D \int_{-\infty}^{\infty} dx \sin^2 \theta \cos \phi \sin \phi \]

\[ = S H_D \int_{\theta(x=\infty)}^{\theta(x=-\infty)} d\theta - p \lambda \sin \theta \cos \phi \sin \phi \]

\[ = 2 S H_D \lambda \sin \phi \cos \phi \]

\[ = S H_D \lambda \sin 2\phi \] (A.40)

The integrals over the first two terms are 0 because \( \frac{\partial \hat{M}}{\partial x} \bigg|_{x=\infty} = \frac{\partial \hat{M}}{\partial x} \bigg|_{x=-\infty} = 0 \) and \( \hat{z} \cdot \frac{\partial \hat{M}}{\partial \phi} = 0 \).
\[ T_4 = \int dV \tau_4 \]
\[ = S \frac{\alpha}{\gamma} \int_{-\infty}^{\infty} dx \left( \frac{\partial \hat{M}}{\partial \phi} \right)^2 \frac{d\phi}{dt} + \left( \frac{\partial \hat{M}}{\partial \theta} \cdot \frac{\partial \hat{M}}{\partial \phi} \right) \frac{\partial \theta}{\partial q} \frac{dq}{dt} \]
\[ = S \frac{\alpha \lambda}{\gamma} \frac{d\phi}{dt} \int_{-\infty}^{\infty} dx \sin^2 \theta \]
\[ = S \frac{2 \alpha \lambda}{\gamma} \frac{d\phi}{dt} \int_{\theta(x = -\infty)}^{\theta(x = \infty)} d\theta - p \sin \theta \]
\[ = 2 S \frac{\alpha \lambda}{\gamma} \frac{d\phi}{dt} \]

Adding all “torques” results in equation A.42 which is identical to the second part of equation 2.11 in the main text.

\[ \sum T = 0 \]
\[ p \frac{dq}{dt} + \alpha \lambda \frac{d\phi}{dt} = \gamma \lambda \pi \left( H_x \cos \phi + H_y \sin \phi \right) \]

(A.42)

### A.3 Adding pinning

To perform depinning calculations, a pinning site and \( z \) field are added to the model, again using the approach of Ryu [44], see page 19. This results in equation A.43.

\[ \frac{dq}{dt} - p \lambda \frac{d\phi}{dt} = F_{pin} - p \gamma \lambda H_z \]
\[ p \frac{dq}{dt} + \alpha \lambda \frac{d\phi}{dt} = \gamma \lambda \pi \left( H_x \cos \phi + H_y \sin \phi \right) \]

\[ F_{pin} \] is the pinning force. We model the the pinning potential as a finite-ranged harmonic potential well. In reality the pinning potential at an anisotropy boundary is not harmonic, using a more realistic approximation potential would be interesting for further research. The maximum pinning force can be calculated from the anisotropy step:

\[ F_{pin,max} = (K_{eff,non-irradiated} - K_{eff,irradiated}) \frac{2 \lambda}{\delta} \tanh \frac{\delta}{2 \lambda} \]

(A.44)

This equation is taken from literature [40] page 17, and \( \delta \) is the width of the transition region form the irradiated to the non-irradiated area, which is approximately 10 nm in our samples.
The equations are not solved analytically by hand, *mathematica* is used for the calculations presented in section 4.3 and 5.3.2.
Appendix B

Preparatory experiments

Before the actual measurements are performed, some sample properties need to be checked. Is the e-beam evaporated golden surface sufficiently smooth to grow an out-of-plane magnetic structure on top? When contacts are protected using pen, can this pen be removed completely or will it leave a thin insulating layer? Is it possible to do wirebonding on the chosen materials? These questions will be answered in this section.

To do this, an unpatterned test sample is fabricated. This sample build up out of layers consisting of the same materials with the same thicknesses as final, patterned samples. First 14 nm of titanium and 140 nm of gold are evaporated on the sample. The corners are protected using black pen and a layer of 10 nm of SiO$_2$ is sputtered. The sample is cleaned with an oxygen plasma for 5 minutes and 5 nm Ta, 4 nm Pt and 0.6 nm Co are deposited. Then on half of the the sample 4.0 nm of Pt is deposited, identical to the patterned samples. On the other half only 1.0 nm of Pt is deposited, as a preparation for possible future experiments. The pen, together with the material deposited on top, is removed using acetone. No pen left overs can be observed using a microscope.

A hysteresis loop is measured using a MOKE setup for both regions with different Pt thicknesses. The results are shown in figure B.1. The remanent magnetization 100%, indicating that the system is out-of-plane. This is the desired behaviour and it indicates that the surface quality of thick golden layer is sufficient. The coercive field is approximately 5 mT, but this will probably be larger for the nanostrips which will contain less nucleation sites, simply because of their smaller size.

Next the wirebonding is tested on the areas that had been covered with pen. This is an easy process compared to previous attempts on samples with a Pt top layer. We expect this is related to the fact that the wirebonding wires consist of gold themselves. On two corners the sample is connected to the chip carrier with two wires. The resistance is now
Appendix B. Preparatory experiments

Figure B.1: Hysteresis loops of an unpatterned test sample with the same material stack as the systems that are investigated in the results section of this report. Both graphs are averages over 5 loops.

measured using a multimeter and turns out to be 1.4 Ω. We aimed for a resistance below 10 Ω, so this suitable our experiments. However, it must be realized that the resistance will increase significantly when a small structure is used instead of a thin film.

These short experiments indicate that the materials and thicknesses we use have the desired properties. We can now proceed to examine structured samples and perform the depinning measurements.
Bibliography


