Tailored Properties of Ferromagnetic Thin Films

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Dissertation presented at Uppsala University to be publicly examined in Häggsalen, Ångström laboratorium, Lägerhyddsvägen 1, Uppsala, Thursday, December 11, 2008 at 10:15 for the degree of Doctor of Philosophy. The examination will be conducted in English.

Abstract

Magnetic thin films and patterned nanostructures have been studied with respect to their magnetic properties using SQUID-magnetometry, magnetic force microscopy, electrical measurements, and micromagnetic calculations.

Properties of vortex domain walls, trapped in Permalloy nanowires with artificial constrictions, were investigated experimentally and by numerical calculations. In particular, the geometrical extent and strength of the pinning potential were evaluated. In these wires, long-range vortex domain wall displacement induced by spin polarized alternating currents was obtained numerically at reduced threshold current densities as compared with the direct current case. Due to the asymmetry of the energy potential, the long-range displacement direction is determined by the vortex chirality.

Strained FeCo/Pt superlattices with strong perpendicular anisotropy were investigated experimentally. The strain was controlled by varying the thickness of each alternating layer with monolayer precision and was found to have a dominating effect on the total anisotropy.

Epitaxial films of the diluted magnetic semiconductor (Ga,Mn)As were studied with focus on how the ferromagnetic transition temperature could be controlled by post-growth annealing. The ferromagnetic transition temperature was enhanced by approximately 85% for a Mn-doping concentration of 6% under certain conditions.

A method to manipulate micrometer sized magnetic particles on patterned arrays of elliptical Permalloy microstructures was studied. Controlled motion and separation of the magnetic particles were obtained using applied rotating magnetic fields. The domain structure of the elliptical elements was studied numerically.

Keywords: micromagnetics, domain wall, vortex, pinning potential, nanowire, spin dynamics, spin transfer torque, magnetic anisotropy, MFM, Permalloy, PMA, magnetic multilayer, Curie temperature, DMS, GaMnAs, magnetic bioseparation

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ISSN 1651-6214
urn:nbn:se:uu:diva-9403 (http://urn.kb.se/resolve?urn=urn:nbn:se:uu:diva-9403)
Till mina föräldrar
This thesis is based on the following papers, which are referred to in the text by their Roman numerals.

I  **Asymmetric energy potential of pinned domain walls at constrictions**  

II  **Long-range vortex domain wall displacement induced by an alternating current: Micromagnetic simulations**  

III  **Non-adiabatic spin transfer torque in high anisotropic magnetic nanowires with narrow domain walls**  

IV  **Magnetic anisotropy of tetragonal FeCo/Pt (001) superlattices**  

V  **Perpendicular magnetocrystalline anisotropy in tetragonally distorted Fe-Co alloys**  

VI  **Influence of annealing parameters on the ferromagnetic properties of optimally passivated (Ga,Mn)As epilayers**  

VII  **Simulations of magnetic microstructure in thin film elements used for programmable motion of magnetic particles**  
VIII  Magnetization of thin-film Permalloy ellipses used for programmable motion of magnetic particles

IX  Domain configuration of Permalloy ellipses in a rotating magnetic field

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<td>Lattice parameter</td>
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<td>Frequency</td>
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<td>Constants</td>
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<td>Non-adiabaticity</td>
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<td>$\delta$</td>
<td>Tip oscillation amplitude</td>
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<td>Cantilever oscillation</td>
<td>$h$</td>
<td>Thickness; Planck’s constant</td>
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<td>Wall width parameter</td>
<td>$h_{\text{lift}}$</td>
<td>Lift height</td>
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<td>Elementary charge</td>
<td>$\eta$</td>
<td>Cantilever damping factor</td>
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<td>Anisotropy energy density</td>
<td>$H$</td>
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<td>Anisotropy energy</td>
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<td>Boltzmann constant</td>
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<td>Exchange energy</td>
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<td>Effective spring constant</td>
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<td>Surface energy</td>
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<td>$\kappa$</td>
<td>Lattice type parameter</td>
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<td>$E_Z$</td>
<td>Zeeman energy</td>
<td>$\chi$</td>
<td>Phase shift</td>
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A multitude of technological applications in modern society rely on magnetism. Generally, magnetic materials are divided in two groups: soft and hard. Soft magnetic materials, characterized by high magnetic permeability and their ability to enhance magnetic flux, are used in devices such as power generators, transformers, inductors, relays, and electromagnets. Hard magnetic materials, or permanent magnets, characterized by their ability to stay magnetized, are used in devices such as power generators, electrical motors, and loud speakers. The applications mentioned above utilize magnetism in bulk materials. If the dimensionality of the system is reduced, a thin film is obtained with magnetic properties different from those of bulk. In this regime, the film thickness becomes an important parameter that controls the magnetization for a given material. Moreover, advances in micro-engineering technology have enabled fabrication of magnetic microstructures and nanostructures with well-defined geometry, which provides additional degrees of freedom to the system.

These structures play an important role in electronic devices and are expected to be indispensable in forthcoming spintronic devices, which exploit not only the charge of the electrons but also their spin. The research on spintronics emerged from the discovery of the giant magnetoresistance effect [1, 2], which already has become implemented for example in the non-volatile magnetoresistive random access memory (MRAM) and in the read-heads of today’s hard disk drives. Prospective spintronic devices include the spin-transfer torque (STT) MRAM and the racetrack memory [3] that employ spin polarized currents to manipulate the magnetization. The evolving field of spintronics would be increasingly fueled if magnetism could be added to semiconductor nanostructures that are compatible with silicon based technology.

This thesis deals with magnetism in thin films. Particularly, how the magnetic properties of films can be tailored for specific applications. Magnetic films and patterned structures have been fabricated and analyzed experimentally. In addition, numerical calculations have been employed to study the static and dynamic behavior of the magnetic configuration. The constituent chapters of this thesis are organized as follows. Chapter 2 describes the methods employed to fabricate the thin films and pattern the surfaces. A recipe is contained that details the steps in the nanostructure fabrication process. Chapter 3 reviews some experimental techniques used to extract the magnet-
ic properties of the thin films. These include techniques for: imaging magnetic domains, measuring the macroscopic magnetic moment of a film, and measuring the magnetoresistance of a single patterned structure. In chapter 4, some aspects of micromagnetism are discussed. Energy and anisotropy terms that constitute the total energy are defined and scrutinized. Spin polarized current as a driving force in the micromagnetic equation of motion is addressed. Chapter 5 is concerned with domain walls in ferromagnetic nanowires. It is shown how domain walls can be nucleated, trapped, and displaced at artificial pinning sites. The energy potential of these pinning sites is calculated and their heights are measured experimentally. Chapter 6 introduces the reader to the field of diluted magnetic semiconductors, where the ferromagnetic transition temperature is still below room temperature for silicon-compatible compounds. Here, we show how the ferromagnetic transition temperature in Mn-doped GaAs compounds can be enhanced by a post-growth annealing treatment. Chapter 7 discusses a method to manipulate micrometer-sized magnetic particles on surfaces with patterned arrays of elliptical Permalloy elements.
2 Sample Fabrication

The thin films used in this work have been fabricated by different deposition techniques, which will be described briefly in this section. The growth process is carried out in a deposition chamber at high vacuum. A source material that is to be deposited (evaporant) is vaporized by some means of energetic excitation. The vapor then condenses on all surfaces in the chamber, including the substrate surface. Under suitable conditions, the evaporated constituents obtain enough energy to travel ballistically towards the exposed surfaces and the sample growth can be controlled by a mechanical shutter. A quartz crystal microbalance (QCM) is exposed to the vapor and used to measure the total thickness and the growth rate.

2.1 The Substrate

The choice of substrate is crucial since its properties heavily influence the properties and quality of the deposited films. If epitaxial crystal growth is intended, the lattice parameters of the substrate must match those of the film material. A lattice mismatch will produce stress in the layers close to the substrate [4]. If the film thickness is sufficiently large this stress can be relaxed by means of dislocations. A large mismatch can also result in a 45° rotation of the grown crystal, which can be seen for example in body centered cubic (bcc) Fe grown on face centered cubic (fcc) MgO [5]. The in-plane lattice parameter $a$ of Fe and MgO is 2.86 Å and 4.21 Å, respectively. This behavior is seen in paper IV and V, where bcc FeCo grows on fcc Pt with a 45° rotation of the in-plane lattice axis. For the fabricated Fe$_{0.36}$Co$_{0.64}$/Pt superlattices the lattice parameter $a$ of Fe$_{0.36}$Co$_{0.64}$ and Pt is 2.65-2.72 Å and 4.1-4.3 Å, respectively, depending on the interlayer thicknesses. Substrates of MgO are used in papers IV and V, GaAs in paper VI, and Si in papers I, III, and IX.

2.2 Thermal Evaporation

In this technique, the evaporant is thermally activated, which results in material meltdown. When a constant temperature and vapor pressure are obtained, the shutter is opened and evaporated material condenses on the sub-
strate. Thermal evaporation has the advantage of being simple to implement and allows a wide range of materials to be deposited. The film uniformity, on the other hand, is negatively affected by the low average kinetic energy of the evaporated atoms impinging on the film surface. Here, two types of thermal evaporation will be mentioned that employ different heating mechanisms.

2.2.1 Resistive Evaporation
Thermal evaporation by resistive means is one of the simplest methods of thin film deposition. A small quantity of the evaporant is placed on a filament, made of tungsten, tantalum, or some other thermally resistant material. The filament is then subjected to an electrical current, which makes the evaporant dissolve into vapor. This deposition method is suitable for materials with low melting points. Impurity levels can be kept down by making sure the filament has a negligible vapor pressure at the working temperature in comparison to the evaporant, but even so the impurity levels are much higher than the other deposition methods described here. The Permalloy elements in paper IX are fabricated using resistive evaporation. The deposition rate is typically in the order of 0.1 to 1 nm/s.

2.2.2 Electron Beam Evaporation
Thermal evaporation can be performed by letting a beam of electrons heat the evaporant, which is placed inside a thermally resistant crucible [6]. In this case, the beam can be accurately focused on the target material, which reduces the contamination from surrounding materials that entails evaporation by resistive heating. Electrons are thermionically emitted from a filament cathode inside an electron gun and accelerated towards the crucible, which also acts as the anode. To avoid condensation of evaporated material on the filament, the electron gun is located away from the crucible so that the electron beam trajectory from the filament to the crucible makes an angle of more than 180 degrees. The electron beam can produce extremely high evaporant temperatures, making possible the deposition of almost any material. The films in paper I are grown using an acceleration voltage of typically 10 kV with stable and low growth rates in the range 0.02 to 0.05 nm/s.

2.3 Molecular Beam Epitaxy
In principle, molecular beam epitaxy (MBE) is based on thermal evaporation although it does not share the simplicity of the evaporation methods discussed above. MBE is used for growing epitaxial films, i.e. monocrystalline films grown relative to the monocrystalline substrates. The growth is per-
formed by directing fine beams of evaporated material from effusion cells towards a substrate. In order to minimize collisions with residual gas and to avoid contamination, the deposition is carried out under ultra-high vacuum (UHV) conditions with a base pressure ~ $10^{-7}$ Pa. The substrate temperature can be regulated during growth and has an important impact on the film properties [7]. Vapor pressures of the evaporants determine the growth rate, which is sometimes kept as low as 1 pm/s, and can be controlled by varying the source temperature of the individual evaporants. This allows atomic layer-by-layer growth with accurate composition and high uniformity. The films in paper VI are fabricated using MBE.

2.4 Magnetron Sputter Deposition

The sputter deposition technique employs electric discharge of an inert sputter gas. At electric field strengths above the breakdown voltage of the gas, the gas atoms are ionized. Positive ions from the ionized gas are accelerated towards the target, which also acts as the cathode and consists of the material to be deposited. For high enough energies of the impacting ions, atoms will break loose from the target, propagate through the plasma, and condense on surrounding surfaces. Electrons liberated in the ionization process will travel towards the anode and may, upon collision with atoms, further ionize the gas. For the plasma to be self-sustained the density of the inert gas must be sufficiently high, corresponding to pressures up to about 10 Pa [8]. The idea of magnetron sputter deposition is to use an applied magnetic field as a catalyst for the ionization [9]. Electrons from the plasma and secondary electrons emitted from the target will get trapped by the magnetic field lines and accelerated towards the anode. Due to the Lorentz force, the electrons will follow helical trajectories. Thus, their total travelled path is increased and also the probability of collisions with gas atoms. Magnetron sputter deposition therefore allows a reduction of the sputter gas pressures by one up to two orders of magnitude. For all samples in this work fabricated with magnetron sputtering argon was used as the sputter gas with a base pressure below $10^{-7}$ Pa. The films in paper IV and V are fabricated using magnetron sputter deposition.

2.5 Nanostructure Patterning

The nanostructures studied in this work were fabricated using electron beam lithography (EBL) and material deposition in a lift-off process. This technique allows the final sample shape to be defined before the material deposition takes place. The procedure used to fabricate the samples is summarized in Figure 2.1 and Table 2.1.
In a first step the substrate surface is covered with photo-resist. This is done by spinning the substrate as the resist is applied. For small features, i.e. in the nanometer regime, a thinner resist is usually required and therefore higher angular speeds are used during spin coating. Large features, such as electrodes, generally require a thicker layer of resist due to the higher beam currents used. Two types of resist are available: positive and negative. A positive resist is weakened and becomes more soluble when exposed by electrons, whereas a negative resist is strengthened and becomes less soluble. A positive resist is convenient in processes where the coverage of the patterned structures is small compared to unoccupied area. Positive PMMA was used in paper IX and positive ZEP 520A was used in paper I. An electron beam is then scanned over the resist, following a predefined pattern. The total exposure time depends on the beam current, the exposed area, and the required charge dose.

The parts of the positive resist that are exposed to the electron beam become weakened and are easily dissolved with a chemical developer. In order to strengthen the remaining resist, the substrate can be heat treated. At a next stage, the material is deposited using a suitable evaporation technique. Unless a mask is used, the deposited material mantles the whole substrate surface. Finally, a lift-off step is performed using chemical solvents that dissolves the remaining resist. Resist is sometimes particularly apt to remain, for example in spaces between closely separated deposited structures. Additional ultrasonic treatment is the standard procedure for these situations. In
the case of nanostructures composed of only one material, a single-step lift-off process is sufficient. If additional features e.g. electrodes of another material are implemented, a multiple-step process is required.

Table 2.1. Recipe for electron beam lithography using ZEP 520A.

1. Cleaning  The substrate is put in acetone for one hour and then subjected to ultrasonic treatment for 3 minutes at 45 Hz. After drying the substrate, it is put in iso-propanol and subjected to ultrasonic treatment for 3 minutes at 45 Hz.

2. Spin coating A small quantity of ZEP 520A is applied on the substrate surface. A mechanical spinner then rotates the substrate at either (i) 5 seconds at 1000 Hz followed by 60 seconds at 3000 Hz for large features, e.g. electrode pads, or (ii) 5 seconds at 1000 Hz followed by 60 seconds at 5000 Hz for small features, e.g. the nanowire parts.

3. Drying  The substrate is baked at 180 °C for 30 minutes.

4. EBL  The desired pattern is written onto the resist using a 50 kV scanning electron microscope. For nanowires a charge dose of 130 μC/cm² and a beam current of 30-50 pA is used, and for electrodes a charge dose of 140 μC/cm² and a beam current of 300-500 pA is used. For structures of intermediate sizes a beam current of 80-150 pA is typically used.

5. Development  After electron beam exposure, the pattern is developed in ZED N50 for 100 seconds and baked at 120 °C for 10 minutes.
3 Characterization Techniques

A number of different techniques were used to study the thin films prepared in this thesis. For characterization of the crystal structures, X-ray reflectivity (XRR) and X-ray diffraction (XRD) were used. For analysis of the chemical composition, Rutherford backscattering (RBS), X-ray photo-electron spectroscopy (XPS), or secondary ion mass spectroscopy (SIMS) was used. The structure of the superlattices in paper IV and V is described in [10] and details of the single crystal film in paper VI can be found in [11]. Scanning electron microscopy (SEM) was routinely used to study the shape and size of the patterned nanostructures. The principles of SEM are thoroughly reviewed in [12]. Below, the techniques used to study the magnetic properties of the films will be discussed.

3.1 SQUID

A superconducting quantum interference device (SQUID) is one of the most sensitive tools for measuring the magnetic moment of a sample [13]. An MPMS-XL SQUID magnetometer from Quantum Design [14] was used to study the magnetic properties of the samples in papers IV and V. The MPMS-XL system uses liquid helium to keep the magnet as well as the SQUID element in a superconducting state. The dewar contains vacuum radiation shields that reduce helium boil off (a helium consumption of slightly more than 5 liters/day is typical). A persistent superconducting current flowing in a closed superconducting circuit, including the superconducting magnet, can maintain highly stable magnetic fields up to 5 T with a field stability of typically 1 ppm/hour. Thin film samples are mounted on a strip of paper within a plastic straw. This straw is mounted on a sample rod connected to a stepper motor that allows precise motion of the sample inside the SQUID pickup coil. Measurements were carried out using the reciprocating sample option (RSO). The RSO option oscillates the sample inside the pickup coil and the magnetic response measured by the SQUID is fitted to the expression of an ideal magnetic dipole moment moving with respect to the pickup coil with the same oscillation as the sample (see Figure 2.2). The oscillation amplitude can be as large as 4 cm and oscillation frequency can be varied from 4 Hz to below 1 Hz. Two different scanning methods are available: the maximum slope and the center scan methods. The former me-
Method oscillates the sample in a region with respect to the pickup coil where its magnetic response has its maximum slope. Since this region is less than one centimeter in length, the measurements can be carried out at a comparably high oscillation frequency which prevents the sample from being subjected to large field variations. The center scan method oscillates the sample over a region of a few centimeters including the peak signal of the magnetic response. During these measurements the sample position is well defined and the recorded magnetic moments are more accurate than with the maximum slope method. On the other hand, these measurements are more time consuming. A sensitivity of $1 \times 10^{-11}$ Am$^2$ in a field of 0.25 T can be achieved. The temperature stability can be controlled within 0.5% in the range 1.9 K to 400 K and the intrinsic magnetic field uniformity is typically 0.01% of the set value over a 4 cm sample scan ($\pm$ 2 cm from the center of the pickup coil).

![Figure 2.2 SQUID response from a ferromagnetic thin film (black line) and the fit of a magnetic dipole expression (gray line). The total oscillation amplitude is 4 cm.](image)

### 3.2 MFM

One way to image the magnetic microstructure is to use magnetic force microscopy (MFM), which is a modification of atomic force microscopy (AFM) [15]. In MFM, as being the case with other scanning probe microscopy (SPM) techniques, a probe is traversed over the sample surface. The
probe consists of a magnetized tip [16] that interacts directly with the sample and a flexible cantilever, which deflections are monitored optically [17]. The measurements were carried out in the tapping/lift mode using the force gradient detection method [18]. The cantilever is scanned twice over the surface. During the first scan, the topography is obtained using the root-mean-square of the cantilever oscillation amplitude as feedback. A rescan is then performed, following the topography at a constant lift height, giving the magnetic contrast from the measured change in resonance frequency. The basics of this technique will be briefly discussed here. More details can be found in [18-20].

3.2.1 Principle of Operation

The cantilever of the scanning probe has a resonance frequency $\omega_0$ that depends on its spring constant $k$ and effective mass $m$ according to

$$\omega_0 = \sqrt{\frac{k}{m}}. \quad (3.1)$$

If the cantilever is driven with sinusoidal oscillations of frequency $\omega$ and amplitude $\delta_0$, the tip will oscillate with the same frequency $\omega$, an amplitude $\delta$, and a phase shift $\chi$. In absence of external forces, and under assumption that the oscillation amplitudes are small compared to the cantilever length, the probe-sample separation can be described by the equation of motion corresponding to a damped harmonic oscillator [19]:

$$m \frac{\partial^2 s}{\partial t^2} + \frac{\omega_0}{Q} \frac{\partial s}{\partial t} + k(s - s_0) = \delta_0 k \cos(\omega t), \quad (3.2)$$

where $s_0$ is the tip-sample distance at zero oscillation amplitude and $s = s(t)$ is the time dependent tip-sample distance under oscillations. The material specific quality factor $Q = m\omega_0/(2\eta)$ includes the damping factor $\eta$. The solution to equation (3.2) is

$$s(t) = s_0 + \delta \cos(\omega t + \chi) \quad (3.3)$$

with a tip oscillation amplitude

$$\delta = \frac{\delta_0 \omega_0^2}{\sqrt{(\omega^2 - \omega_0^2)^2 + 4\eta^2 \omega^2}} \quad (3.4)$$
and phase shift

\[
\chi = \arctan \frac{2\eta \omega}{\omega^2 - \omega_0^2}.
\]

(3.5)

If an external force \( \mathbf{F} \) is present, the motion can be described by adding the \( z \) component of this force \( F_z \) on the left-hand side of (3.2). Assuming that the oscillation amplitudes are small compared to the tip-sample distance and considering a first order Taylor series of \( F_z = F_z(s, \partial s/\partial t) \), the spring constant is shifted by the vertical force gradient [19],

\[
k_{\text{eff}} = k - \frac{\partial F_z}{\partial z}.
\]

(3.6)

As a result, the cantilever resonant frequency will be modified according to

\[
\omega' = \omega_0 \sqrt{1 - \frac{1}{k} \frac{\partial F_z}{\partial z}}.
\]

(3.7)

Looking at the shift in resonance frequency \( \Delta \omega \) we have

\[
\Delta \omega_0 = \frac{\omega'_0 - \omega_0}{\omega_0} = -\frac{1}{2k} \frac{\partial F_z}{\partial z}
\]

(3.8)

if expression (3.7) is expanded to first order in \( 1/k \cdot \partial F_z/\partial z \). From (3.4) and (3.5) we now see that force gradients encountered by the probe during a scan can be obtained experimentally as a change in oscillation amplitude \( s \) and/or phase shift \( \chi \).

Concerning the sensitivity of the force-gradient method, limitations originate from thermal noise from the cantilever, noise from the displacement sensor, and noise from the oscillation control amplifier [18]. The minimum detectable force gradient is given by

\[
\left( \frac{\partial F_z}{\partial z} \right)_{\text{min}} = \frac{1}{\delta_{\text{rms}}} \sqrt{\frac{2 f_{\text{bw}} k_B T}{\omega_0 Q}},
\]

(3.9)

where \( \delta_{\text{rms}} \) is the root mean square of the cantilever oscillations and \( f_{\text{bw}} \) is the measurement bandwidth. The Q-factor \( Q \) varies for different materials but is also dependent on the surrounding medium and pressure and ranges from below 40 for some gases up to over 100000 for UHV [21, 22]. Minimum detectable force gradients can be in the \( \mu \text{N/m} \) range for 1 kHz bandwidth,
corresponding to forces in the order of fN [23]. It should be noted that only a high $Q$ does not necessarily provide optimal measurement conditions. The settling time $\tau$ for a cantilever to reach a steady state oscillation can be written $\tau = 2/Q\omega_0$ [18]. Thus, a high Q-factor of e.g. 50000 and a resonance frequency of 50 kHz would give a bandwidth of 0.5 Hz, which is impractical for most applications.

3.2.2 Tip-Sample Interaction

The stray field from each magnetic moment in a sample will add to form a collective stray field $\mathbf{H}$ according to [24]

$$\mathbf{H}(r) = \frac{1}{4\pi} \left( \int_{S''} \mathbf{n}_s \cdot \mathbf{M}_1(r'') \frac{\mathbf{R}}{|\mathbf{R}|} dS'' - \int_{V''} \nabla \cdot \mathbf{M}_1(r'') \frac{\mathbf{R}}{|\mathbf{R}|} dV'' \right),$$  

(3.10)

where $\mathbf{n}_s$ is the out-of-plane normal from the sample surface, $\mathbf{M}_1(r'')$ is the magnetization in the sample at $r''$, and $\mathbf{R} = r - r''$. The first integral is taken over the sample surface $S''$ and the second over the sample volume $V''$. The terms $\mathbf{n}_s \cdot \mathbf{M}_1$ and $-\nabla \cdot \mathbf{M}_1$ correspond to the surface magnetic charge density and volume magnetic charge density, respectively. As the magnetized tip is scanned over the sample it will interact with this stray field and experience a force $\mathbf{F}$ according to [24]

$$\mathbf{F}(r) = \mu_0 \int_{V'} \nabla (\mathbf{M}_2(r') \cdot \mathbf{H}(r' + r)) dV',$$

(3.11)

where the integration is carried out over the tip volume $V'$. Here, $\mathbf{M}_2(r')$ denotes the magnetization in the tip at $r'$. If we consider only the force component $F_n$ along the normal $\mathbf{n}$ to the cantilever we get the force gradient along the same direction as

$$\frac{dF_n}{dn} = \mathbf{n} \cdot \nabla (\mathbf{n} \cdot \mathbf{F}) = \mu_0 \int_{V'} \mathbf{n} \cdot \nabla (\mathbf{n} \cdot \nabla (\mathbf{M}_2(r') \cdot \mathbf{H}(r' + r))) dV'.$$

(3.12)

If the normal $\mathbf{n}$ coincides with the vertical component $z$ and the tip is approximated by a magnetic point dipole moment $\mu$ with $\mu_x$, $\mu_y$, and $\mu_z$ being its respective $x$, $y$, and $z$ components, expression (3.12) can be written as

$$\frac{dF_n}{dz} = \mu_0 \left( \mu_x \frac{\partial^2 H_x}{\partial z^2} + \mu_y \frac{\partial^2 H_y}{\partial z^2} + \mu_z \frac{\partial^2 H_z}{\partial z^2} \right).$$

(3.13)
In this expression, $H_x$, $H_y$, and $H_z$ are the $x$, $y$, and $z$ components of the sample stray field at the position of the tip, respectively. For a tip magnetized along the $z$ direction, this expression reduces to

$$\frac{dF_n}{dz} = \mu_0 \mu_z \frac{\partial^2 H_z}{\partial z^2}. \quad (3.14)$$

If the tip-sample distance is large compared with the length of a contributing dipole moment, as is typically the case under the measurement mode discussed here, the stray field from an arbitrary magnetic dipole $\mu$ can be approximated by [25]

$$H(r) = \frac{1}{4\pi} \left( \frac{3(r \cdot \mu) r}{r^5} - \frac{\mu}{r^3} \right). \quad (3.15)$$

By employing (3.14) and (3.15), the expected response of a magnetic force microscope from a magnetic sample can be calculated (see Figure 5.18).

### 3.3 Electrical Measurements

Electric currents can be employed to probe some properties of magnetic nanostructures. It has been shown that vortex walls and transverse walls can be distinguished by their anisotropic magnetoresistance [26]. In fact, even the chirality of each of these domain wall types can be determined by the same method. By using the giant magnetoresistance (GMR) effect [1, 2] the domain wall position in a nanowire can be measured [27]. A magnetic domain wall trapped in a nanostructure can be excited to resonance by using a time dependent driving force. The resonance frequency is dependent on the structure geometry and can be determined by electrical measurements [28, 29] as we will see below.

#### 3.3.1 Power Calibration

If the impedance of the sample does not match the impedance of the transmission lines, microwave reflections will occur at the boundary. This impedance mismatch creates a strong frequency dependence of the generator output power. The absorbed power increases the sample temperature and hence the sample resistance measurements at different frequencies become inaccurate due to the strong correlation between resistance and temperature. In Figure 3.1a the sample resistance is shown as a function of the injected power level. In order to feed a constant power to the sample, a calibration procedure is necessary. The resistance is measured as a function of frequency for
different power levels and a power correction for each frequency is determined and stored for the measurement. In Figure 3.1b the resistance of a sample is measured with, and without, power calibration.

Figure 3.1 Data obtained from electric measurements of Permalloy discs. (a) Sample resistance as a function of the injected power level. (b) Resistance versus injected microwave frequency with (squares), and without (circles), power calibration. The power level is 16 dBm. (c) Magnetoresistance recorded with an applied magnetic field swept parallel (0 deg) and perpendicular (90 deg) to the electric current, respectively. (d) Measured $U_{DC}$ for a disc with a diameter of 1 μm and a thickness of (40 ± 10) nm. The crossings with zero voltage occur at 284 MHz and 322 MHz.
Figure 3.2 Evolution of a magnetic vortex state in a Permalloy disc as an applied magnetic field $H_A$ is increased in the direction of the arrow. The dark and bright regions correspond to local magnetic moments (arrows) aligned parallel and perpendicular to the field, respectively. (I) As no field is applied the vortex core sits in the centre of the disc. (II) As the field is increased ($H_A = 0.8H_s$) the vortex core moves perpendicular to the field direction. (III) The vortex core reaches the rim and annihilates. This transient state is shown for $H_A = H_s$. (IV) The magnetic structure finally relaxes ($H_A = H_s$) into a single domain state with all moments aligned parallel to the applied field.

3.3.2 Sample Magnetoresistance

A current $j$ passing a region with a gradient in the magnetization $M$ experiences an anisotropic magnetoresistivity

$$\rho(\theta) = \rho_0 + \Delta\rho \cos^2(\theta)$$

(3.16)

that depends on the angle $\theta$ between the current and the magnetization vector [30]. The parameter $\rho_0$ denotes the resistivity for a magnetization orthogonal to the current and $\Delta\rho = \rho_1 - \rho_0$ where $\rho_1$ is the resistivity for a magnetization parallel to the current. If an external field is applied with an angle to the electric current, the effect from Hall resistivity must be added to (3.16) due to the Lorentz force. Figure 3.1c shows a typical magnetoresistance graph for a Permalloy disc. The magnetic microstructure of the disc changes as an external magnetic field is applied. The expected behavior has been studied in micromagnetic simulations. Figure 3.2 shows a magnetic vortex state in a
disc with a diameter of 1 µm and thickness of 40 nm. As the applied field \( H_A \) is increased the vortex core propagates perpendicular to the applied field direction towards the rim. For a sufficiently large field \( H_A = H_s \) it annihilates and all local magnetic moments in the disc become aligned along the applied field. In this state, the anisotropic magnetoresistance gives rise to a resistance splitting for measurement currents parallel and perpendicular to the spins. This splitting can be seen in Figure 3.1c and amounts to 90 mΩ.

3.3.3 Homodyne Detection

One way to study the resonance frequency of a magnetic microstructure is to use a homodyne detection method [28]. If the magnetic configuration is oscillating with a frequency \( \omega \), its change in magnetoresistance will have a frequency dependence according to

\[ \Delta R(\omega) = A_1(\omega)\sin(\omega t + \chi) + A_2(\omega)\sin(2\omega t + \zeta) + \ldots \]  

(3.17)

If the magnetic vortex is excited by the injection of a spin polarized alternating current \( I = I_0\sin(\omega t) \), the measured voltage change can be written as \( U = I\Delta R \). By using the trigonometric sum rules we see that this voltage consists of two parts, one that is time dependent and one that is time independent \( U_{DC} \). The latter can be written

\[ U_{DC} = \frac{I_0A_1(\omega)}{2}\cos(\chi), \]  

(3.18)

where \( \chi \) is the phase shift between the current and the first frequency component. At resonance, the phase lag of a weakly damped oscillator is \( \chi = \pi/2 \) and the measured voltage \( U_{DC} \) will be zero. In Figure 3.1d the measured \( U_{DC} \) is shown for a Permalloy disc with a diameter of 1 µm and thickness of (40 ± 10) nm. The first crossing with zero voltage occurs at a frequency of 284 MHz. As a comparison, a resonance frequency of 272 MHz has been reported for a Permalloy disc with a diameter of 1.1 µm and thickness of 40 nm obtained from vector network analyzer resonance spectra [31].

The experimental setup used for the electrical measurements is illustrated in Figure 3.3. A lock-in amplifier SRR830 with a measuring current of 6 kHz is used to probe the resistance between ports A and B. A microwave source MG3692B is connected to the circuit and allows the injection of signals with a well defined frequency in the sub-GHz range and an output power up to about 23 dBm. The gray circular object in the figure symbolizes a Permalloy disc and the shaded rectangular objects illustrate gold electrodes that partly overlap the disc. Capacitors are blocking the low frequency signals and inductors are blocking the high frequency signals.
Figure 3.3 The experimental setup for the electrical measurements. In the bottom right corner a SEM image shows a Permalloy disc with a diameter of 2 μm.

3.4 XMCD-PEEM

Magnetic domain imaging for some of the samples (including the ones in paper III) was carried out using X-ray magnetic circular dichroism photo emission electron microscopy (XMCD-PEEM) with synchrotron radiation. The radiation is produced in the synchrotron by undulators and guided into beam lines. By tuning the beam energy it is possible to extract element specific information from the atomic absorption lines [32]. The magnetic contrast is obtained by subtracting data from X-rays of different polarizations. Another advantage with synchrotron X-rays compared with smaller lab-friendly sources is the wide beam with almost parallel radiation, which allows a large sample surface to be illuminated simultaneously. In this section a brief overview of the technique will be given together with a description of the used experimental setup. More details can be found in [28].
Figure 3.4 Photoemission spectra of epitaxial Fe on GaAs(001). The graphs for right and left circularly polarized photons are symbolized with a solid and dashed line, respectively, and their difference with a dotted line. The L3 peak corresponds to transitions $2p_{3/2} \rightarrow 3d$ and the L2 peak corresponds to transitions $2p_{1/2} \rightarrow 3d$. Adapted from [33].

The mechanism of X-ray magnetic circular dichroism (XMCD) relies on the different absorption coefficients for left and right circularly polarized radiation. In Figure 3.4 the photoemission spectrum is shown for iron. The magnitude of the absorption peaks, representing the L2 and L3 edges, respectively, depends on the helicity of the X-rays. Electrons from the 2p core level states are excited to empty 3d levels. Due to the spin orbit coupling the 2p states are split into the $2p_{1/2}$ and $2p_{3/2}$ states. The transition $2p_{1/2} \rightarrow 3d$ corresponds to the L2 peak and the transition $2p_{3/2} \rightarrow 3d$ corresponds to the L3 peak. The background of the spectrum originates from transitions into empty s and p states. In ferromagnets, the density of states for 3d electrons at the Fermi level is spin dependent. Therefore the X-ray absorption intensities are different for different polarizations of the light. In our measurements, data are collected for both helicities. The magnetization is then obtained by taking the difference between spectra of opposite helicities. The spin and orbital moments can be determined using the XMCD sum rules. However, since the main interest was to image magnetic domains these sum rules [34, 35] will not be reviewed here.

Photo emission electron microscopy (PEEM) is a common technique for surface imaging [36]. In the used experimental setup, the grounded sample is
mounted inside the PEEM column in front of the PEEM extractor lens. Electrons released from the sample surface are accelerated by 10 keV to an objective lens. Subsequently the electrons pass an octupole stigmator lens, where corrections for aberrations can be made, and two projective lenses, where the field of view can be controlled. A microchannel plate amplifies the number of electrons that finally impinge on a luminescent screen. Images are taken with a CCD camera. Here, the majority of released, and analyzed, electrons are not photoelectrons but secondary electrons released in an Auger process [32].

**Figure 3.5** Experimental setup at BESSY. (A) Sample load-lock. (B) Sputter gun. (C) Main vacuum chamber. (D) Sample transfer pipe. (E) Vacuum pump. (F) Spin detector. (G) Sample intermediate chamber. (H) CCD camera. (I) PEEM column. (J) Vacuum pump. (K) Channel plate. (L) Beam line.

Using the Fe absorption edge we achieved nice contrast and a magnetic resolution of about 30 nm. Typically, a set of 10 images for each helicity were taken and averaged to improve the signal to noise ratio. The XMCD-PEEM measurements in this work were carried out at *Berliner Elektronenspeicherring-Gesellschaft für Synchrotronstrahlung* (BESSY). A photo of the experimental setup is shown in Figure 3.5.
4 Micromagnetism

Micromagnetic theory describes the magnetic microstructure and magnetization dynamics of ferromagnetic bodies on length scales too large for a quantum mechanical treatment. This phenomenological approach started with a publication by Landau and Lifshitz [37]. Inspired by Bloch's earlier work [38] they investigated the spin distribution in regions between opposite magnetizations, i.e. magnetic domain walls. Rigorous development of the theory was made by Brown [39-41]. The idea is to substitute the individual spins in a small volume of the material with uniform spin directions by a magnetization vector. The extent to which this substitution is valid can be determined by properties of the material, where an upper limit is given by its exchange length $l_{ex}$. Under the influence of an effective field which has contributions from various magnetic energies, the magnetization in the volume changes direction. Neighboring magnetization vectors are allowed to vary only by a small angle, thus giving rise to a continuous vector field [41]. The dynamics of these vectors can be calculated using the Landau-Lifshitz equation of motion [37].

Micromagnetic theory can be used on a wide range of system sizes. The lower limit of applicability corresponds to length scales where spins no longer behave continuously due to their quantum properties. A theoretical upper limit does not exist. In practice, when implementing micromagnetic theory in numerical calculations an upper limit originates from the computational capacity.

The numerical calculations in papers I, III, and VII-IX were performed using the object oriented micromagnetic framework (OOMMF) [42]. The numerical calculations in paper II were performed using a code developed by Nakatani [43].

4.1 Energetics

Micromagnetic theory is based on the minimization of the total energy. The magnetization directions $\mathbf{m}(\mathbf{r}) = \mathbf{M}(\mathbf{r})/M_s$ are chosen so that the total energy of the system is at a global or local minimum under the constraint $\mathbf{m}^2 = 1$. In the following sections, the different contributions to the total energy will be discussed.
4.1.1 Zeeman Energy

A magnetic moment will try to reduce its energy by aligning itself parallel to an applied magnetic field \[44\]. The energy that describes the interaction of a magnetic moment with an applied field \( \mathbf{H}_A \) is called the Zeeman energy \( E_Z \):

\[
E_Z = -\mu_0 \int V \mathbf{M} \cdot \mathbf{H}_A dV. \tag{4.1}
\]

4.1.2 Exchange Energy

The exchange energy \( E_{ex} \) is due to the exchange interaction between magnetic spins. If we consider two neighboring spins with spin vectors \( \mathbf{S}_1 \) and \( \mathbf{S}_2 \), respectively, the exchange energy can be written

\[
E_{ex} = -2J \mathbf{S}_1 \cdot \mathbf{S}_2, \tag{4.2}
\]

where \( J \) is the exchange coupling term \[45\]. The sign of \( J \) determines the type of coupling between spins. If \( J < 0 \), anti-parallel alignment is preferred and the coupling is referred to as antiferromagnetic. For \( J > 0 \), parallel alignment is favored and the coupling is ferromagnetic. Since neighboring spins are only allowed to vary by a small angle \( \phi \) we can approximate \((4.2)\) by a Taylor series. If we furthermore neglect the first term, which does not depend on \( \phi \), the exchange energy between two spins \( i \) and \( j \) can be written

\[
E_{ex} = -2JS^2 \cos \phi = JS^2 \phi^2 = JS^2 \left| \mathbf{s}_j - \mathbf{s}_i \right|^2 = JS^2 \left( (\mathbf{e}_j \cdot \nabla)\mathbf{s} \right)^2 \tag{4.3}
\]

with \( \mathbf{e}_j \) symbolizing the position vector from spin \( i \) to spin \( j \) and \( \mathbf{s}_j = \mathbf{S}_j / |\mathbf{S}_j| \) [41]. The exchange energy of an ensemble of spins is thus given by

\[
E_{ex} = \frac{1}{2} JS^2 \sum_{nn} \left( (\mathbf{e}_j \cdot \nabla)\mathbf{s} \right)^2, \tag{4.4}
\]

where the summation is taken over nearest neighbors. In the continuous limit, \((4.4)\) can be expressed as \[46\]

\[
E_{ex} = A \int \left( (\nabla S_x)^2 + (\nabla S_y)^2 + (\nabla S_z)^2 \right) dV. \tag{4.5}
\]

The exchange stiffness constant \( A = JS^2 \kappa / a \) depends on the lattice parameter \( a \) and the crystal structure with \( \kappa = 1 \) for a simple cubic, \( \kappa = 2 \) for a body centered cubic, and \( \kappa = 4 \) for a face centered cubic lattice.
4.1.3 Demagnetization Energy

The demagnetization energy (or magnetostatic self-energy) comes from the interaction of a magnetic body with its own stray field. Because of its dependence on the body geometry, this energy is sometimes also called shape anisotropy energy. The stray field $H_d$ can be defined using Maxwell’s equation

$$\nabla \cdot B = \mu_0 \nabla \cdot (H + M) = 0,$$

(4.6)

where $B$ is the magnetic flux density [47]. This gives, if we consider $H_d = H$ to be the field produced by the divergence of the magnetization vector [44],

$$\nabla \cdot H_d = -\nabla \cdot M.$$

(4.7)

In magnetized ferromagnetic bodies, the free poles appearing on the ends will generate a magnetic field oriented antiparallel to the magnetization. A relation between the demagnetizing field and magnetization can be written

$$H_d = -N M,$$

(4.8)

where the demagnetizing factor $N$ depends on the body geometry [48]. In, for example, an infinite cylinder with its long axis parallel to the $z$ axis, $N = 0$ for vectors parallel to $z$ and $N = \frac{1}{2}$ for vectors perpendicular to $z$. In a sphere $N_x = N_y = N_z = \frac{1}{3}$. More generally, $N$ takes the form of a tensor. The demagnetization energy $E_d$ of a body is found by integration over its volume

$$E_d = -\frac{1}{2} \mu_0 \int_V H_d M dV.$$

(4.9)

As can be seen, the minus signs in (4.8) and (4.9) cancel out and give rise to a positive energy. The equations above give a full description of the stray field energy. Equivalent results can be obtained by looking at the individual contributions from magnetic dipole moments in the body. These moments produce a stray field according to (3.10) and can be used in expression (4.9) to find the total stray field energy.

4.2 Magnetic Anisotropy

If the internal energy of a body is independent on the spontaneous magnetization direction, the body is magnetically isotropic [49]. However, this case is a rare exception and it turns out that most magnetic materials are instead
magnetically anisotropic. The strength, or degree, of the anisotropy is measured as an energy, which is called the magnetic anisotropy energy. A distinction is made between anisotropies related to the crystal structure and anisotropies due to deviations from the case of a perfect crystal lattice. Shape anisotropy, originating from the sample geometry, is discussed in section 4.1.3. Here, different types and origins of magnetic anisotropy will be discussed in terms of energy densities $e_{an}$. The total anisotropy energy $E_{an}$ is found by integration over the corresponding body dimensions.

4.2.1 Magnetocrystalline Anisotropy

The magnetocrystalline anisotropy arises as a result of coupling between the spin and orbital motion of electrons in the crystal lattice. This gives rise to stable directions of the spontaneous magnetization, or easy axes, along which the magnetization preferably aligns. Because of its origin, the magnetocrystalline anisotropy is strongly related to the lattice symmetry. In order to find the directional dependence of the magnetization, expansion of the anisotropy energy density in terms of the direction cosines of the magnetization vector is often used [50], where the expansion is such that it fulfills the symmetry requirements set by the crystal lattice. The direction cosines, i.e. the components along the $x$, $y$, and $z$ directions of the magnetization unit vector $m$ are usually denoted $m_1$, $m_2$, and $m_3$, respectively. Only the first two non-vanishing terms of the expansion will be considered, since higher order terms usually give negligible contributions to $e_{an}$. Here, we will look at two cases of magnetocrystalline anisotropy.

(i) Uniaxial Anisotropy

The simplest case of magnetocrystalline anisotropy is uniaxial anisotropy. This type of anisotropy is manifested in hexagonal and tetragonal structures. As an example, we can look at hexagonal cobalt where the easy axis is parallel to the crystal $c$ axis, having a polar angle $\theta = 0$. As the magnetization is rotated away from this axis, the anisotropy energy increases and reaches a maximum for a polar angle $\theta = 90^\circ$, after which it decreases to a second minimum at $\theta = 180^\circ$. The anisotropy energy density is found by expansion in terms of the direction cosine $m_3 = \cos(\theta)$:

$$e_{an} = -K_{u1} m_3^2 - K_{u2} m_3^4,$$

where we have limited the expansion at the fourth order term. The anisotropy constants $K_{u1}$, and $K_{u2}$, measured in units of J/m$^3$, are both positive in cobalt [51]. A large positive $K_{u1}$ describes an easy axis, whereas a large negative $K_{u1}$ describes an easy plane. Figure 4.1 shows the anisotropy energy calculated using (4.10) for different choices anisotropy constants.
Figure 4.1 Uniaxial anisotropies represented by energy surfaces. The length of the plotted radial component is proportional to the energy density for that direction. The anisotropy constants are chosen to illustrate different cases at similar energy scales. (a) Easy perpendicular direction with $K_{u1} = 1 \text{ J/m}^3$, $K_{u2} = 0 \text{ J/m}^3$. (b) Easy plane with $K_{u1} = -0.9 \text{ J/m}^3$, $K_{u2} = -0.9 \text{ J/m}^3$. (c) Easy cone, $K_{u1} = -2 \text{ J/m}^3$, $K_{u2} = 0 \text{ J/m}^3$. (d) Easy perpendicular and easy plane with $K_{u1} = -4.8 \text{ J/m}^3$, $K_{u2} = 4.8 \text{ J/m}^3$.

(ii) Cubic Anisotropy
The magnetocrystalline anisotropy energy density of cubic crystals can be found by a polynomial series expansion in terms of the directional cosines according to [44]

$$e_{\text{an}} = K_{c1} (m_1^2 m_2^2 + m_2^2 m_3^2 + m_3^2 m_1^2) + K_{c2} (m_1^2 m_2^2 m_3^2), \quad (4.11)$$

where $K_{c1}$ and $K_{c2}$ are anisotropy constants. If $K_{c1} > 0$, the $<100>$ directions are the easy axes. If $K_{c1} < 0$, the anisotropy depends on the sign and magnitude of $K_{c2}$. Figure 4.2 shows the anisotropy energy calculated using (4.11) for different choices of anisotropy constants.
Figure 4.2 Cubic anisotropies represented by energy surfaces. The length of the plotted radial component is proportional to the energy density for that direction. The anisotropy constants are chosen to illustrate different cases at similar energy scales. (a) Easy <100> directions with, $K_{c1} = 1.82 \text{ J/m}^3, K_{c2} = 0 \text{ J/m}^3$. (b) Easy <110> directions with $K_{c1} = -2 \text{ J/m}^3, K_{c2} = 8 \text{ J/m}^3$. (c) Easy <111> directions with $K_{c1} = -1.82 \text{ J/m}^3, K_{c2} = 0 \text{ J/m}^3$. (d) Hard <111> directions with $K_{c1} = 0 \text{ J/m}^3, K_{c2} = 10.94 \text{ J/m}^3$.

4.2.2 Surface Anisotropy

The magnetic surface anisotropy of a ferromagnet results from the reduced symmetry at the surface. A body with cubic crystal structure will have a magnetic surface anisotropy energy density $e_s$ described by

$$e_s = \frac{1}{2} K_{s1} (m_1^2 n_1^2 + m_2^2 n_2^2 + m_3^2 n_3^2) + K_{s2} (m_1 m_2 n_1 n_2 + m_2 m_3 n_2 n_3 + m_3 m_1 n_3 n_1),$$

(4.12)

where $\mathbf{n} = (n_1, n_2, n_3)$ is the surface normal [41]. $K_{s1}$ and $K_{s2}$ are anisotropy constants measured in units of J/m$^2$. If $K_{s1} = K_{s2}$, as is the case of a structurally isotropic medium, the energy density (4.12) reduces to

$$e_s = \frac{1}{2} K_s (\mathbf{n} \cdot \mathbf{m})^2.$$

(4.13)
A positive $K_s$ will minimize $E_s$ for a magnetization vector perpendicular to the surface plane, and gives $e_s = K_s$ when magnetization lies in plane.

### 4.2.3 Interface Anisotropy

Similarly to the magnetic surface anisotropy, this type of anisotropy is also a consequence of symmetry reduction at the boundary. The two cases can be distinguished by observing that the environment of a surface is of non-solid state character, whereas an interface is a boundary between two materials. A structure with a large number of interfaces is called a multilayer. If we consider a multilayer where the magnetic layers are interspaced with non-magnetic layers, the effective anisotropy of each magnetic layer can in the simplest case be expressed as

$$K_{\text{layer}} = K_v + 2 \frac{K_s}{h}$$

(4.14)

with volume- and surface-anisotropy constants, $K_v$ and $K_s$, respectively. The surface contribution is inversely proportional to the layer thickness $h$ of the magnetic material and consequently becomes large in very thin layers [44].

### 4.2.4 Magnetoelastic Anisotropy

When a ferromagnet is magnetized, it expands (or contracts) in the direction of the magnetization and hence its shape changes. This phenomenon is called magnetostriction. The opposite case, when stress induces a change in the magnetization, is referred to as magnetoelasticity. A typical expansion of the magnetoelastic anisotropy in cubic crystals results in the energy density

$$e_{\text{me}} = \frac{3}{2} \lambda_{100} \sum_i \sigma_{ii} m_i^2 - 3 \lambda_{111} \sum_{i>k} \sigma_{ik} m_i m_k,$$

(4.15)

where higher order terms are neglected [52]. Here, $\lambda_{100}$ and $\lambda_{111}$ are material specific magnetostriction coefficients describing the strength of the magnetoelastic interaction, and $\sigma_{ik}$ are the components of the stress tensor. This stress is considered to be of non-magnetic origin and can be external mechanical stress, internal stress created by heterogeneity in structure or composition, or it can be due to lattice imperfections such as lattice dislocations creating local internal stress. Stress of magnetic origin is treated, e.g. in [44]. If we consider a structure, which is elastically isotropic, and uniaxial stress that makes an angle $\theta$ with the magnetization direction, the magnetoelastic anisotropy energy density reduces to
\[ e_{\text{me}} = -\frac{3}{2} \lambda \sigma \cos^2 \theta. \]  

(4.16)

The stress \( \sigma = \varepsilon \varepsilon \) is related to the strain \( \varepsilon \) via the elastic modulus \( \varepsilon_e \).

4.2.5 Induced Anisotropy

The induced magnetic anisotropy describes the influence of deviations from perfect crystallinity. These deviations can be generated by certain treatment procedures: annealing [51]; magnetic annealing, in which an external magnetic field is applied during the heat treatment [51, 53, 54]; applied magnetic [55] or electric [56] field during the fabrication process. Magnetic anisotropy can also be induced by lattice misfit [57] between a thin film and the substrate used in the deposition process and is referred to as strain-induced anisotropy. Strain-induced anisotropy is studied in more detail in papers IV and V.

4.3 The Effective Field

Summing up the energy contributions from external field, exchange coupling, demagnetization, and anisotropy energy we can write the total free energy as

\[ E_{\text{tot}} = E_Z + E_{\text{ex}} + E_d + E_{\text{an}}. \]  

(4.17)

The effective field \( H_{\text{eff}} \) interacting on \( M \) is obtained by differentiating the total energy with respect to the magnetization [58]

\[ H_{\text{eff}} = -\frac{1}{\mu_0 V} \frac{\partial E_{\text{tot}}}{\partial M}. \]  

(4.18)

As Brown pointed out [41], the equilibrium magnetization configurations can be found at the minima of this energy by solving the variational problem

\[ \frac{\partial E_{\text{tot}}}{\partial \mathbf{m}} = 0. \]  

(4.19)

This leads to the equilibrium condition that the magnetization vectors at all points should align parallel to the local effective field and that the normal derivatives of the magnetization components should be zero at the surface:
\[ \mathbf{m} \times \mathbf{H}_{\text{eff}} = 0, \quad (4.20a) \]
\[ \mathbf{m} \times \frac{\partial \mathbf{m}}{\partial \mathbf{n}} = 0. \quad (4.20b) \]

In the boundary condition above, zero surface anisotropy has been assumed. If surface anisotropy is present, a term \( \frac{\partial \mathbf{e}}{\partial \mathbf{m}}(2A)^{-1} \) must be added on the left hand side of (4.20b) [44]. The conditions (4.20a) and (4.20b) allow the equilibrium magnetization configurations of a system to be found. The difficulty of finding exact solutions can be overcome by introducing a convergence criterion on (4.20a) and use iterative methods to find a numerical solution to the micromagnetic configuration. However, such an approach only provides static solutions to the problem. If one wishes to study the magnetization dynamics, the magnetic equation of motion described in sections 4.5-4.9 have to be used.

### 4.4 The Exchange Length

One of the important length scales in micromagnetism is the exchange length \( l_{\text{ex}} \), which is the distance over which the magnetization is expected to change direction in the presence of magnetic anisotropies [59]. This length occurs where the exchange energy is in equilibrium with the dominating anisotropy term. When the shape anisotropy is dominant, the exchange length is given by

\[ l_{\text{ex}} = \sqrt{\frac{A}{K_d}}, \quad (4.21) \]

where \( A \) is the exchange stiffness constant and \( K_d = \mu_0 M_s^2 / 2 \) is the stray field energy constant. Using standard values for Permalloy, \( A = 13 \text{ pJ/m} \) and \( M_s = 860 \text{ kA/m} \), the exchange length amounts to \( l_{\text{ex}} = 5.3 \text{ nm} \).

### 4.5 Landau-Lifshitz Equation

Starting with the origin of atomic magnetism, we recall that the magnetic moment \( \mathbf{\mu} \) of electrons obeys a linear relationship with the angular momentum \( \mathbf{J} \) according to

\[ \mathbf{\mu} = -\gamma_0 \mathbf{J}, \quad (4.22) \]
where the gyromagnetic ratio $\gamma_0 = \mu_0 ge/(2m_e)$ is dependent on the Landé g-factor $g$, the absolute value of the electron charge $e$, and the electron mass $m_e$ [41]. The torque due to a magnetic field is found from the time derivative of the angular momentum

$$\frac{d\mathbf{J}}{dt} = \mathbf{\mu} \times \mathbf{H}. \quad (4.23)$$

This describes the precession of the magnetic moment about a magnetic field. The angular frequency $\omega$ of this precession is proportional to the field strength by $\omega = \gamma_0 H$ [41]. Using (4.22) and (4.23) together with the definition of the magnetization $\mathbf{M} dV = \sum \mathbf{\mu}_i$, where the summation is taken over all atomic magnetic moments $\mathbf{\mu}_i$ in $dV$, we get for the volume element $dV$

$$\frac{d\mathbf{M}}{dt} = -\gamma_0 \mathbf{M} \times \mathbf{H}. \quad (4.24)$$

The model by Landau and Lifshitz is based on this equation and describes the dynamics of a magnetization vector in an effective field $\mathbf{H} = \mathbf{H}_{\text{eff}}$, which takes into account contributions from all involved energy terms $E_{\text{tot}}$ according to (4.18). Equation (4.24) describes the precession of a magnetization vector that is constant over time. If the system is subjected to dissipative processes, we expect the time derivative of $\mathbf{M}$ to approach zero after some time so that we have a static situation with the magnetization pointing along $\mathbf{H}_{\text{eff}}$. A term describing this dissipation and that pushes the magnetization vector towards $\mathbf{H}_{\text{eff}}$ can be added to (4.24), giving the Landau-Lifshitz (LL) equation

$$\frac{d\mathbf{M}}{dt} = -\gamma_0 \mathbf{M} \times \mathbf{H}_{\text{eff}} - \frac{\lambda}{M_s} \mathbf{M} \times (\mathbf{M} \times \mathbf{H}_{\text{eff}}). \quad (4.25)$$

The second term on the right is referred to as the Landau-Lifshitz damping term and contains a material specific constant $\lambda > 0$.

### 4.6 Landau-Lifshitz-Gilbert Equation

The micromagnetic equation of motion was derived by Gilbert in 1955 using a different approach. He observed that (4.25) could be obtained from a Lagrangian formulation using the magnetization components $M_x$, $M_y$, and $M_z$ as generalized coordinates [60]. Effects of dissipation is then introduced as a viscous force, with components proportional to the time derivate of the magnetization, giving the Landau-Lifshitz-Gilbert (LLG) equation
\[
\frac{dM}{dt} = -\gamma_G M \times H_{eff} - \frac{\alpha}{M_s} M \times \frac{dM}{dt},
\]
(4.26)

where two new parameters are introduced: \(\gamma_G\) and \(\alpha\). The two equations, (4.25) and (4.26), are mathematically similar, which becomes clear if we rewrite (4.26) on the following form

\[
\frac{dM}{dt} = -\frac{\gamma_G}{1 + \alpha^2} M \times H_{eff} - \frac{\gamma_G \alpha}{(1 + \alpha^2)M_s} M \times (M \times H_{eff}).
\]
(4.27)

In fact, we see that (4.25) and (4.26) are identical with the following choices of the parameters \(\gamma_0\) and \(\lambda\) in (4.25):

\[
\gamma_0 = \frac{\gamma_G}{1 + \alpha^2},
\]
(4.28a)

\[
\lambda = \frac{\gamma_G \alpha}{1 + \alpha^2}.
\]
(4.28b)

The question now arises concerning the physical difference between the two equations of motion, (4.25) and (4.26). For small damping there is usually no significant difference and taking the limit of zero damping they both express the same physics. However, taking the limit of infinite damping we get

\[
\lim_{\alpha \to \infty} \frac{dM}{dt} = 0
\]
(4.29a)

for (4.26), and

\[
\lim_{\lambda \to \infty} \frac{dM}{dt} = \infty
\]
(4.29b)

for (4.25). In other words, in the case of large damping, expression (4.29a) suggests a slow loss of energy and slow approach to the low-energy state, whereas expression (4.29b) predicts a rapid loss of energy where the system quickly reaches its low-energy state. This observation may suggest that the Gilbert type of damping gives a physically more realistic description. However, both damping terms are being extensively used and the validity of the two models is still under debate [60-62].
4.7 The Damping Parameter

The damping parameter $\alpha$ becomes important when the magnetization dynamics is to be scrutinized. To illustrate this, the relaxation of a magnetization vector is studied by means of micromagnetic simulations. The Landau-Lifshitz equation of motion is used on a single cubic cell with 5 nm side length. Initially the magnetization is directed along a vector that makes an azimuthal angle of 30° with the $x$ axis and a polar angle of 45° with the $z$ axis. At $t = 0$, an external field $\mu_0 H_A = -50$ mT is applied along $z$. As shown in Figure 4.3 the system shows no tendency of damping for $\alpha = 0.000001$ and the magnetization keeps precessing around the $z$ axis. Increasing $\alpha$ one order of magnitude results in a barely distinguishable oscillatory magnetization relaxation towards the applied field direction (not shown). For $\alpha = 0.01$ the system shows a damped behavior and relaxes in 95 ns. Increasing $\alpha$ by one order of magnitude decreases the relaxation time by roughly one order of magnitude. For $\alpha = 0.5$ the oscillations are suppressed and the magnetization vector relaxes in 1.5 ns.

Experimentally, the damping parameter can be estimated e.g. by ferromagnetic resonance measurements [63], time resolved measurements [64], or electrical measurements [65]. The damping parameter is material specific [63] and has been measured for Permalloy, yielding estimates between 0.013 and 0.03 [64-66]. In this thesis, $\alpha = 0.02$ is generally used in simulations of Permalloy. In some cases, when the static magnetic configuration is sought for and the choice of $\alpha$ does not noticeably influence the result, a higher $\alpha$ is used. In papers VII and VIII for example, a value of $\alpha = 0.5$ is used.
Figure 4.3 The relaxation of a magnetization vector for different values of $\alpha$ obtained by micromagnetic simulations based on Landau-Lifshitz equation. The data represents the direction of the magnetization that has its base at the origin of the coordinate axes. The initial state of the vector is the same for the different cases. The displayed time is measured from the application of an external magnetic field (corresponding to a flux density of 50 mT) directed along the negative $z$ direction. A saturation magnetization of $M_s = 860$ kA/m is used.
Figure 4.4 A comparison between the Landau-Lifshitz (LL) and Landau-Lifshitz-Gilbert (LLG) equation of motion for $\alpha = 0.05$, 0.02, and 0.5 respectively. Each figure shows the relaxation of a magnetization vector initially aligned parallel with the $x$ axis and subjected to an applied field $\mu_0 H_A = 5$ mT along the negative $z$ axis at $t = 0$. The left column shows normalized magnetization component in the $x$ direction, whereas the right column shows the normalized magnetization along the $z$ direction. The phase separation of the time dependent magnetization between the LL and LLG cases is indistinguishable for $\alpha = 0.05$ but becomes apparent for $\alpha = 0.2$ and larger.

4.8 Gilbert Damping versus Landau-Lifshitz Damping

In order to evaluate the two different damping mechanisms of the Landau-Lifshitz and Landau-Lifshitz-Gilbert equations, micromagnetic 3D simulations are performed for a Permalloy system using the same initial conditions. A cube with the geometry $10 \times 10 \times 10$ nm$^3$ is defined using a computational mesh of $5 \times 5 \times 5$ nm$^3$ cells. The microstructure is first saturated along $x$, which is parallel to a side of the cube. Then, an external magnetic field $\mu_0 H_z = -5$ mT is applied. The magnetization response is shown in Figure 4.4 for different choices of the damping parameter $\alpha$. Upon application of $H_z$ the
magnetization vector enters a state of damped oscillations about the equilibrium position of least energy. For \( \alpha \leq 0.05 \) there is essentially no difference between Landau-Lifshitz (equation (4.25)) and Landau-Lifshitz-Gilbert (equation (4.26)) damping. For \( \alpha \sim 0.2 \) a phase shift becomes apparent between the two mechanisms. The phase shift in \( M_e \) increases with time and amounts to 0.7 ns lag of LLG after 15 ns. There is no difference in \( M_e \) amplitude between the two mechanisms, not even for \( \alpha = 0.5 \). At \( \alpha = 0.5 \), the damping is strong enough to inhibit repeated oscillations and the system relaxes in 10 ns. The phase lag of \( M_e \) for LLG increases with \( \alpha \) and amounts to 1 ns after 5 ns.

4.9 Spin-Transfer Effects

The idea that electric currents can have an effect on the magnetization was first explored by Berger [67-69]. He argued that the spins of itinerant electrons couple to the local spins of the system through the exchange interaction and thereby exchange spin angular momenta. A spin polarized current passing local spins, with a perpendicular spin component with respect to the spin direction of the current, will experience a change in the angular momentum direction. Consequently, an opposite torque will be exerted on the local spins. This hints at a possibility to control regions of deviating spin alignment, for example domain walls, with spin polarized currents.

4.9.1 Adiabatic Spin-Transfer

The spin transfer effect is implemented in the equation of motion by a spin transfer term \( \tau \) [70],

\[
\tau = -(u \cdot \nabla)M,
\]

(4.30)

where the generalized velocity

\[
u = \frac{gP\mu_B}{eM_s} j
\]

(4.31)

is proportional to the current density \( j \) and the spin polarization \( P \). Introducing the torque \( \tau \) into equation (4.26) gives

\[
\frac{dM}{dt} = -\gamma_0 M \times H_{eff} + \frac{\alpha}{M_s} M \times \frac{dM}{dt} - (u \cdot \nabla)M.
\]

(4.32)
This equation describes the case where a flow of electron spins adiabatically follow the local magnetization which is typically the situation for wide walls. Another form of transfer term \( \tau_z \) has been proposed [71] that describes similar dynamics as long as the length of the magnetization vector \( \mathbf{M} \) is constant:

\[
\tau_z = -\frac{1}{M_s^2} \mathbf{M} \times (\mathbf{M} \times (\mathbf{u} \cdot \nabla)\mathbf{M}). \tag{4.33}
\]

Current induced domain wall motion has been studied by Thiaville and co-workers in [72] using equation (4.32). Continuous domain wall motion could be obtained for a combination of applied field and current, with increasing domain wall velocity up to a critical breakdown field (the Walker field) [73]. Current induced motion without a field could be obtained above a threshold current density, which exceeded the experimental values by one order of magnitude. The torque (4.33) was studied in [71] and the simulations showed that injection of a current pushed the domain wall a distance in the order of 0.1 \( \mu \text{m} \) before it halted. However, when the current was removed the domain wall went back to its initial position. Neither of these results accurately describe experimental observations.

![Figure 4.5](image)

**Figure 4.5** Velocities \( v \) of a transverse domain wall in a 120 nm wide and 5 nm thick nanowire obtained from micromagnetic simulations based on the Landau-Lifshitz-Gilbert equation. Graphs are plotted for different values of \( \beta \) in a perfect wire (a) and in a wire with edge roughness (b). The average grain size at the edges is described by the parameter \( D \). The generalized velocity \( u \) is proportional to the spin polarized current density. Open symbols correspond to nucleation of vortex walls. Experimentally, \( u \) is available in the range indicated by the shaded area. The dashed lines are linear fits with an offset of 25 m/s. Adapted from [70].

### 4.9.2 Non-Adiabatic Spin-Transfer

The torque in the previous section is introduced with the assumption that the spins of the itinerant electrons are always aligned with the local magnetic
moments. To correct for this assumption, an additional non-adiabatic term is introduced to (4.32) giving [70],

\[
\frac{dM}{dt} = -\gamma M \times H_{\text{eff}} + \frac{\alpha}{M_s} M \times \frac{dM}{dt} - (\mathbf{u} \cdot \nabla)M + \frac{\beta}{M_s} M \times ((\mathbf{u} \cdot \nabla)M),
\]

(4.34)

where \(\beta\) is the non-adiabaticity parameter. The non-adiabatic spin-transfer effect becomes apparent in thin walls. The threshold current density becomes zero in simulations using equation (4.34) for perfect wires and a value of \(\beta\) between 0.01 and 0.1 (Figure 4.5). However, in a wire with artificial edge roughness the threshold current density is non-zero. This indicates the influence of the extrinsic pinning sites. Simulated velocities of induced domain wall motion are about two orders of magnitude larger than experimental values, which may be due to stronger pinning effects in experimental systems. A \(\beta\) value of about 0.04 was obtained by experimental measurements of the domain wall velocity \(v\) [70, 74, 75]. Further corrections to (4.34), including the effect of extrinsic pinning and its influence on the threshold current, have been addressed in [76]. A non-adiabatic term as an extension of (4.33) has been proposed in [77]. This term is on the same form as the non-adiabatic term in (4.34) and yields a corresponding \(\frac{\beta}{\gamma} = \frac{\xi}{1+\xi^2}\) where \(\xi = \frac{t_{\text{ex}}}{t_{\text{sf}}}\) is the ratio between the exchange time \(t_{\text{ex}}\) and the spin flip time \(t_{\text{sf}}\).

### 4.9.3 The Spin Polarization

The strength of the spin transfer torque is proportional to the spin current density \(jP\), where \(j\) is the absolute value of \(j\) in (4.31). Thus, the spin polarization \(P\) is an important parameter, which is desired to be as high as possible. By using time-resolved magnetic transmission X-ray microscopy (MTXCM) the spin polarization in Permalloy has been estimated by Kasai and coworkers to be \(P = 0.67 \pm 0.16\) [78].
5 Properties of Magnetic Domain Walls

In a magnetic domain all spins are aligned in the same direction. If all spins in a body are aligned in one direction, it is said to be saturated and its magnetic configuration is referred to as a single domain state. Normally, for a ferromagnetic material in zero field, the spins are not all aligned in one direction unless its size is very small. Instead, the spins are mutually aligned only for separate regions, or domains, and its magnetic structure is in a multidomain state.

The idea of magnetic domains was first introduced by Weiss in 1907 [79]. His molecular field theory described the spontaneous magnetization observed in ferromagnetic materials. Domains were introduced as an explanation to the observed weakly magnetized, or demagnetized, states. Experimental results that supported the existence of domains were first provided by Barkhausen 1919 [80]. He wound a coil around a ferromagnetic sample and connected it to a loudspeaker. When the sample was gradually magnetized by an approaching permanent magnet, he heard a crackling noise. This noise (Barkhausen noise) was explained by voltage pulses generated in the coil from discontinuous steps (Barkhausen jumps) in the magnetization process. The first observation of domain patterns was made by Bitter in 1931 [81]. On a polished surface of a ferromagnetic sample he placed colloidal ferromagnetic particles. By using an optical microscope he could indirectly observe the domains from the distribution and alignment of the colloidal particles. The region separating two domains is referred to as a domain wall. Its properties were first investigated by Bloch 1932, who found that the wall must extend over some distance due to the exchange interaction that opposes abrupt changes in neighboring spin directions [38].

Since the first experimental observations, several sophisticated methods have been developed for observing magnetic domains, for example different variants of transmission electron microscopy (TEM), scanning electron microscopy (SEM), magnetic force microscopy (MFM), neutron topography, magneto-optical techniques, and X-ray spectroscopy. These techniques and others are reviewed and compared in [44]. In this thesis, mostly MFM was used to study the magnetic domains. However, some magnetic domain structures have been studied using XMCD-PEEM.
5.1 The Origin of Domains

Since the formation of domains in a body depends on the magnetic anisotropies and the sample geometry it is not possible to ascribe a single mechanism responsible for the existence of domains [44]. Qualitatively stated, magnetic domains form in order to reduce the energy of a system and stable domain states could be determined by minimization of the total energy. However, a stable domain configuration is not necessarily the global minimum of a system as we will see in section 5.3.1.

Figure 5.1 The region of spins, represented by arrows, between two domains with antiparallel spin configurations constitutes a magnetic domain wall. The type of domain wall is defined by the way their spins rotate. In a Bloch wall (a) the spins rotate in the plane of the domain wall, whereas in a Néel wall (b) the spins rotate perpendicular to the plane of the domain wall.

5.2 Different Types of Domain Walls

In a domain wall the spins gradually change direction going from one domain to the other. Depending on the material properties, specimen geometry, and involved energy terms different spin configurations can be present. Usually, magnetic domains are classified into two categories: 180° walls that separate two domains with opposite magnetizations, and 90° walls that separate domains whose magnetizations are orthogonal. Figure 5.15a shows a magnetic microstructure with five domains and six domain walls. The four domain walls enclosing the central domain are 90° walls, whereas the other two are 180° walls.

Depending on the spin rotation across the wall, two types of domain walls can be distinguished. In a Bloch wall, illustrated in Figure 5.1a, the spins rotate in the domain wall plane. This type is common in bulk samples. In a Néel wall (Figure 5.1b), the spins rotate in a plane perpendicular to the domain wall plane. For thin films, this type of wall has less magnetostatic energy compared with Bloch walls [50]. This energy relation can be understood by considering the magnetic surface charges appearing where a Bloch intersects the film surface: the energy generated by the surface charge in-
creases with decreasing film thickness. Since the corresponding energy in case of a Néel wall is due to volume charges appearing inside the film, its contribution to the domain wall energy density decreases with decreasing film thickness.

By choosing a ferromagnetic material with negligible uniaxial anisotropy and a small film thickness we can expect the magnetization to lie in plane and domain walls to be of Néel type. Then, by restricting the film width to a few hundreds of nanometers we obtain a nanowire (nanostrip) where spins are forced to align along the wire because of the large shape anisotropy. Under these conditions, the transverse wall and the vortex wall can be observed. These magnetic configurations are characterized by their axial magnetization alignment, which can be either head-to-head or tail-to-tail depending on if the magnetizations are pointing towards or away from the wall, respectively. A vortex wall is identified by two additional parameters; the vortex polarity, which is positive for an out-of-plane vortex core magnetization and negative for a magnetization pointing into plane, and the vortex chirality, which can be clockwise or counter clockwise depending on the sense of spin rotation about the vortex core.

Figure 5.2 A phase diagram of domain wall structures obtained from micromagnetic simulations by comparing the energy of three domain wall types in Permalloy nanowires. The dotted line is a fit where the parameter $A$ is 4.99 nm. Adapted from [82].
5.2.1 Domain Wall Type Phase Diagram

The wall energy depends on the type of wall and on the sample geometry. By comparing the energy of transverse and vortex walls, respectively, a phase diagram for domain walls in nanowires can be established [82, 83]. Figure 5.2 shows the regions of least energy, obtained from micromagnetic simulations for the different types of walls including also the asymmetric transverse type of wall [82]. Thus, qualitatively stated, a wide and thick wire favors vortex walls, whereas a thin and narrow wire favors transverse walls. The stability region presented in Figure 5.2 is valid for a perfect wire of constant width and no edge roughness. In a wire with periodically varying width the corresponding phase diagram is expected to be different. For example, in a perfect wire the asymmetric transverse wall is stable in a range of thicknesses between 3 and 8 nm. In a wire with periodical constrictions, the stability region for this type of wall covers a larger parameter space for the same wire width as we will see in the following sections.

![Diagram of curved nanowire with periodical constrictions](image)

**Figure 5.3** A curved nanowire with periodical constrictions with a close-up of the central segment. The nanowire geometry is defined by the thickness $h$, width $w$, constriction depth $d$, and constriction wavelength $l$. 
5.3 Nucleation of Domain Walls

In a straight Permalloy nanowire the shape anisotropy forces the spins to align parallel to the wire. Hence, the remanent magnetic structure is likely to be in a single domain state. To control the domain wall nucleation, external magnetic fields were used with the following procedure. First, the structure is saturated in-plane along its hard axis perpendicular to the long axis of the wire. This field is then gradually decreased to zero. As the switching field is reached a phase transition occurs and the single domain splits in two domains, separated by a domain wall. If the saturating field is applied in the positive $y$ direction (see Figure 5.3), a tail-to-tail domain wall will form upon relaxation of the domain configuration. Conversely, for a saturating field applied in the negative $y$ direction, a head-to-head domain wall will nucleate. Curved nanowires, such as the ones depicted in Figure 5.4 and Figure 5.5 allow for precise control of the domain wall position using the magnetic fields generated by a vector magnet. Another way of nucleating domain walls in a nanowire is to use nucleation pads [84, 85] or current generated Oersted fields [86].

![Figure 5.4 SEM image of a curved Permalloy nanowire with thickness $h = 20$ nm, width $w = 480$ nm, constriction depth $d = 120$ nm, and constriction wavelength $l = 720$ nm. Gold electrodes are connected to the nanowire.](image)

Figure 5.5 XMCD-PEEM image of a nanowire with the same geometry as in Figure 5.4. The magnetic contrast, i.e. the dark and bright regions, reveals two domains of uniform magnetizations. A vortex domain wall separates the two domains and is situated between two of the gold electrodes. The field of view is 15 μm.

5.3.1 Field Sequence Dependent Nucleation

The type of domain wall nucleated in a wire depends on the confining geometry. In addition, the applied field sequence was found to have an influence on the nucleation process. Figure 5.6 shows the evolution of the magnetization configuration under demagnetization of a wire segment with constrictions ($w = 240$ nm, $l = 480$ nm, $h = 20$ nm, and $d = 60$ nm). An in-plane field $H_A$ is used to align all local magnetic moments parallel to the hard axis (Figure 5.6a). As the field is decreased the magnetization configuration remains symmetric with respect to the central hard axis (Figure 5.6b) in a certain field interval. Further reduction of the field will gradually break the symmetry (Figure 5.6c) and a transverse type of domain wall can be discerned. Two distinct patterns are observed according to which the symmetry breaking can occur: the configuration depicted in Figure 5.6c and the configuration obtained by reflection of the magnetization in the same figure with respect to the central hard axis. This symmetry breaking pattern determines the chiral-
ty of vortex walls nucleated by further reduction of the field. In Figure 5.6d, the field has been reduced to $\mu_0 H_A = 20$ mT. A transverse wall, spanned between the two edges of the wire, is clearly visible. None of these configurations show any out-of-plane component. From here, the relaxation is studied for two different field sequences. For field sequence 1, the field is reduced to zero in one step. For field sequence 2, the field is reduced to 0 in steps of 2 mT. Figure 5.7a-d shows the temporal evolution of the in-plane and out-of-plane spin configuration for field sequence 1. At $t = 0$, no out-of-plane component is present (not shown). Due to the abrupt change in Zeeman energy, the local magnetic moments are perturbed and the ones at the notch-free edge obtain a small out-of-plane component. This component increases in magnitude with time and extends through the part of the wall that has an in-plane component aligned parallel with the hard axis (Figure 5.7a,e). The perpendicular component is then released from the edge (Figure 5.7b,f) and manifests itself as the core of a vortex wall. The core propagates (Figure 5.7c,g) into the pinning site confined by the constrictions and pulls the wall along the easy axis of the wire. Due to its gyrotropic nature, the vortex core makes an orbital motion about the pinning site [87] that is strongly correlated with the wall oscillations back and forth along the easy axis. Finally, the nucleated vortex enters an equilibrium state (Figure 5.7d,h). For field sequence 2, the magnetization perturbation due to the stepwise decrease in Zeeman energy is too small to build up any out-of-plane component. The wall contracts with decreasing field but otherwise retains its transverse shape similar to that of Figure 5.6d. At zero field the transverse wall is at stable equilibrium.

The longitudinal magnetization component $M_x$ is illustrated in Figure 5.8a. For both field sequences, the magnetization oscillates as the field is decreased. However, the transverse wall is more strongly damped and oscillates with much reduced amplitudes as compared to the vortex wall. The vortex wall and transverse wall settle with a final normalized longitudinal magnetization of 0.071 and 0.17, respectively. These magnetization offsets indicate the asymmetry of the system that results from the notches. For the vortex, this corresponds to a vortex core offset of 60 nm with respect to the central hard axis. Interestingly, the total relaxation times for the two cases are similar (90 ns and 84 ns for field sequence 1 and 2, respectively). In Figure 5.8b the time evolution of the system energy density $e_{tot}$, calculated as the sum of the exchange and demagnetization energy densities, is illustrated. An abrupt decrease in the energy density is seen at $t = 0$ for field sequence 1 as a result of the nucleation of a magnetic flux closure state. These results indicate the importance of the field history in the nucleation process. Using the same field conditions but reducing the nanowire width produced vortex walls with field sequence 1 and transverse walls with field sequence 2 down to a critical width of 230 nm, after which a transverse wall was obtained for both field sequences.
Figure 5.6 Magnetization configuration in a nanowire with artificial constrictions. In (a), a field $\mu_0 H_A = 1$ T is applied along the $-y$ direction, which is parallel to the local moments represented by the arrows. In (b), the field has been reduced to $\mu_0 H_A = 100$ mT. The magnetization configuration is still symmetric about the central hard axis. In (c), the symmetry is broken for $\mu_0 H_A = 60$ mT. In (d) a transverse domain wall is visible for $\mu_0 H_A = 20$ mT.

Figure 5.7 Time evolution of the domain configuration for field sequence 1. At $t = 0$ the field $H_A$ is abruptly reduced from $\mu_0 H_A = 20$ mT to 0 mT. The left column shows the in-plane magnetization configuration of the non-equilibrium states at $t = 80$ ps (a), 140 ps (b), 500 ps (c), and the equilibrium state at $t = \infty$ (d). The right column shows for the central segment the corresponding magnetization component perpendicular to plane at $t = 80$ ps (e), 140 ps (f), 500 ps (g), and the equilibrium state at $t = \infty$ (h), where the dark and bright regions illustrate magnetization components pointing into the plane and out of the plane, respectively.
Figure 5.8 The dynamics of the relaxation process for field sequence 1 and 2. (a) Time evolution of the normalized magnetization. The initial value of $M/M_s$ is 0.13 and the final values are 0.071 and 0.17 for field sequence 1 and 2, respectively. (b) Time evolution of the energy density $e_{\text{tot}}$. The initial value of $e_{\text{tot}} = 10.1 \text{ kJ/m}^3$ can be compared with the final values 7.1 and 8.3 kJ/m$^3$ for field sequence 1 and 2, respectively.

5.4 Pinning of Domain Walls

The coercivity of a body is often controlled by defects, such as inclusions, voids, crystallographic imperfections, grain boundaries, and surface roughness [88]. These defects generate energy barriers manifested as pinning sites where domain walls can be trapped. In some cases, for example in case of magnetization switching [89-91] or magnetic domain wall transport [27, 92, 93], it may be desirable to have as weak pinning sites as possible since the energy barrier must be overcome by external driving forces. From an application-point-of-view it is also crucial to be able to decrease these driving forces since they are linked to power consumption. Nevertheless, if pinning sites are absent in a sample, the domain wall position becomes sensitive to external force perturbations, which can become a problem in some special cases. Particularly, this problem is apparent in experiments where domain wall motion or nucleation is induced by applied magnetic fields. Likewise, in simulations the absence of pinning sites can lead to unwanted wall propagation when applied fields are used to initialize a magnetic configuration.
In order to have well defined positions where the domain wall can locate itself it is sometimes desirable to create artificial pinning sites. This reduces the number of equilibrium sites for the domain wall and thus enables controlled positioning of the wall between pinning sites. Here, artificial pinning sites in the form of constrictions, or notches, were studied. These are defined by periodically varying the width of a nanowire. A segment of such a wire is shown in Figure 5.3. A domain wall close to a constriction will feel an attractive potential and will be pinned if unaffected by stronger forces. The pinning strength as well as the extent of the pinning site depends on its geometry. A direct measure of the pinning strength is the propagation field $H_p$ required to expel the domain wall from the pinning site.

In the following, the pinning strength of an artificial constriction is evaluated. A vortex domain wall is nucleated and trapped at a pinning site corresponding to Figure 5.7d. The applied flux density $B_x = \mu_0 H_x$ is increased along the long axis of the nanowire in steps of 0.5 mT and for each field the total energy of the system is calculated as the sum of the exchange and demagnetization energies. The result can be seen in Figure 5.9. Reversing either the vortex chirality or the magnetization alignment, between head-to-head and tail-to-tail, reverses the sign of the displacement in the energy potential. Reversing both properties produces the same graphs. The positive (negative) depinning flux densities $B_p$ are 5.25 (-6.25), 6.75 (-9.75), 9.25 (-14.25), 10.25 (-20.25) mT for $d = 20, 40, 60, 80$ nm, respectively. This can be compared with the experimental positive (negative) $B_p$ values on a vortex
with clockwise chirality giving 4.9 (-6.4), 6.6 (-8.8), 8.0 (-9.8), 9.3 (-11.3) mT for $d = 20, 40, 60, 80$ nm, respectively. A vortex with a counter clockwise chirality produced similar positive (negative) $B_p$ values of 5.2 (-6.8), 6.6 (-8.7), 7.5 (-10.2), 8.4 (-11.7) mT for $d = 20, 40, 60, 80$ nm, respectively. The height of the energy barrier in Figure 5.9 is 1.9 (2.4), 2.3 (3.4), 4.7 (8.6), 5.0 (10.2) aJ for the easy (hard) depinning direction and $d = 20, 40, 60, 80$ nm, respectively. These values clearly indicate the asymmetry present in the system. This asymmetry increases with increasing notch depth. A wire with no constrictions shows no asymmetry and the depinning fields are equal for both directions. Similarly, the energy potential of wires with identical constrictions on both edges is symmetric.

Figure 5.10 Average vortex domain wall width for a nanowire with constriction depths $d = 40$ nm and $d = 60$ nm as a function of the applied magnetic flux density $B_x = \mu_0 H_x$. The largest contraction for $d = 40$ nm is 86% for positive fields and 95% for negative fields. The largest contraction for $d = 60$ nm is 81% for positive fields and 90% for negative fields.

A magnetic field applied along the $x$ direction pushes the vortex core along the positive or negative $x$ direction depending on the axial spin alignment. The width of the pinning site is estimated from the total $x$ displacement of the vortex core between the energy maxima seen in Figure 5.9. For $d = 20, 40, 60,$ and 80 nm the width of the pinning site is 66, 212, 210, and 178 nm, respectively. It should be noted that there is also propagation of the vortex core in the transverse direction. The magnitude of this displacement component depends on the constriction geometry and is discussed in paper I. Furthermore, as the domain wall is removed from its energy minimum it con-
tracts with respect to the applied field direction as can be seen in Figure 5.10, where the average domain wall width for $d = 40$ and 60 nm is illustrated. Here, the domain wall width $w_{dw}$ is calculated using

$$w_{dw} = \int_{-\infty}^{\infty} \cos \varphi(x) \, dx,$$

(5.1)

where $\varphi$ is the magnetization angle with respect to the magnetization direction in the domains [44]. Correspondingly, the volume of the domain wall decreases. In the case of $d = 60$ nm, the domain wall volume at zero fields is $7.0 \times 10^{-22}$ m$^3$ but decreases to $6.2 \times 10^{-22}$ m$^3$ for $B_x = -14$ mT and $5.0 \times 10^{-22}$ m$^3$ for $B_x = 9$ mT.

Figure 5.11 Equilibrium state for a 500 nm $\times$ 500 nm cobalt film with a thickness of 0.6 nm. The arrows show the direction of the magnetization, pointing into-plane on the left side and out-of-plane on the right. The magnetic domains are separated by a transition region where the spins gradually rotate 180° in the plane of the domain wall.

5.5 Materials with Uniaxial Anisotropy

In a material with uniaxial anisotropy, the magnetization may be forced to align perpendicular to the film plane. In such a case, head-to-head or tail-to-tail domain walls will not be observed. Instead, the domain walls will attain one of the types illustrated in Figure 5.1. Cobalt is a ferromagnetic material
with strong uniaxial anisotropy. The effective $K_{u1}$ increases with decreasing thickness (see section 4.2.3) according to

$$K_{u1} = -\frac{\mu_0 M_s^2}{2} + K_v + \frac{2K_s}{h}, \quad (5.2)$$

where the first term on the right hand side corresponds to the shape anisotropy, the second term is a volume contribution, and the third term is a surface contribution depending on the thickness $h$ of the cobalt layer. The anisotropy constants $K_v$ and $K_s$ in cobalt have been measured to be $(0.60 \pm 0.07) \text{ MJ/m}^3$ and $(0.58 \pm 0.1) \text{ mJ/m}^2$ at room temperature, respectively [94]. In this thesis, the magnetic microstructure of thin cobalt films is studied. Micromagnetic calculations are performed using an exchange stiffness $A = 16 \text{ pJ/m}$, saturation magnetization $M_s = 1.4 \text{ MA/m}$, damping parameter $\alpha = 0.02$, and a uniaxial anisotropy constant $K_{u1} = 1.5 \text{ MJ/m}^3$. The equilibrium state is calculated in a single cobalt film using a cell size of $2 \text{ nm} \times 2 \text{ nm} \times 0.6 \text{ nm}$, which is small compared to the exchange length in cobalt of $l_{ex} = 3.6 \text{ nm}$. The $x$ and $y$ directions are defined in the same way as in Figure 5.14 and the film spans the volume $0 \leq x \leq 500 \text{ nm}$, $0 \leq y \leq 500 \text{ nm}$, $0 \leq z \leq 0.6 \text{ nm}$. The result is shown in Figure 5.11 with magnetic domains aligned perpendicular to the film plane and separated by a Bloch domain wall. The domain wall profile can be seen in Figure 5.12, which shows the $x$-dependence of the perpendicular magnetization component for a constant $y$ in the cobalt film. Using a hyperbolic tangent [95],

$$f(x) = \tanh\left(\frac{x}{\Delta}\right), \quad (5.3)$$

the domain wall width $w_{dw} = \pi \Delta$ can be estimated from the wall width parameter $\Delta$ yielding $w_{dw} = 22.0 \text{ nm}$. Using (5.1) to calculate the domain wall width gives $w_{dw} = 22.2 \text{ nm}$. In Figure 5.13 the domain wall width is calculated for different $y$ using (5.1). The wall width is essentially constant for $50 \leq y \leq 450 \text{ nm}$ and decreases with less than 15% at the sample edges.
Figure 5.12 The domain wall profile of a Bloch wall in a thin cobalt film. The normalized perpendicular magnetization component is shown (circles) for a constant $y = 249$ nm together with a hyperbolic tangent (straight line) fit to the data.

Figure 5.13 Domain wall width across a thin cobalt film. The average wall width is 22.1 nm.

In a cobalt multilayer, individual cobalt films can couple through the layers and the resulting magnetic microstructure is expected to be different as compared with a single cobalt layer. Micromagnetic calculations are carried out on a cobalt-platinum multilayer Pt 1.8/Co 0.6/Pt 1.2/Co 0.6/Pt 1.2/Co 0.6/Pt
1.8 nm with a lateral size of 500 nm × 500 nm using a computational mesh of 4 nm × 4 nm × 0.6 nm. In the simulations, Pt was assumed to be nonmagnetic. The equilibrium state is shown in Figure 5.14. The magnetization vectors in the different cobalt layers are more or less parallel for constant \( x \) and \( y \). However, this does not apply for the domain walls. In the bottom layer, the magnetization vectors rotate clockwise with respect to the \( y \) axis for increasing \( x \) and the domain wall is of Néel type. In the middle layer, the magnetization vectors rotate clockwise with respect to the \( x \) axis for increasing \( x \) and the domain wall is of Bloch type. In the top layer, the magnetization vectors have a reversed sense of rotation as compared to the bottom layer and the wall is of Néel type. Using (5.3), the domain wall widths for the bottom, middle, and top layers become 21.2, 18, and 21.2 nm, respectively. Equation (5.1) gives domain wall widths of 24.7, 19.6, and 24.7 nm for the bottom, middle, and top layers, respectively.

**Figure 5.14** Magnetization distribution in the different layers of a Co/Pt multilayer.

### 5.6 Manipulation of Domains and Domain Walls

Magnetic domains and domain walls can be controlled using applied magnetic fields as an external driving force [27, 43, 84, 86, 96]. In Figure 5.15
the idea of field driven domain wall motion is illustrated. Figure 5.15a shows a demagnetized five-domain state. As an external magnetic field $H_A$ is applied (see Figure 5.15b), the magnetic domains parallel to the field direction grow at expense of the others. As a consequence, the domain walls separating the domains propagate. In this example, the 180° domain walls propagate perpendicular to the field direction. In the case of head-to-head and tail-to-tail configurations, the domain walls will propagate parallel or antiparallel to the field direction. An example of how a magnetic head-to-head domain wall can be controlled by an applied magnetic field can be found in section 5.4.

Another way of controlling magnetic domains and domain walls is to use spin polarized direct currents (DC) [85, 89, 93, 97, 98]. The driving force of this mechanism is the transfer of angular momentum from the spin polarized electrons to the local magnetic moments [68, 99]. Critical current densities for domain wall motion are in the order of $10^{11}$ - $10^{12}$ A/m$^2$ [74, 85, 93]. For the potential use in applications, it is of interest to reduce this high threshold value. Domain wall displacement by alternating currents (AC) has been predicted below the DC threshold [100]. In paper II, micromagnetic simulations of vortex domain wall motion induced by spin polarized AC is studied.

Consider a head-to-head vortex wall trapped at a pinning site in a nanowire with artificial constrictions. An applied current will push the domain wall in the direction of the electron motion. Thus, an AC will push the domain wall back and forth along the long axis of the wire. The motion of the vortex core on the other hand can be separated from vortex wall motion. Due to its gyrotropic nature the vortex core begins to orbit about the pinning site. If its resonance frequency matches the frequency of the applied current, the
orbital motion of the vortex core is resonantly amplified and could, if the current density is sufficiently high, extend beyond the width of the pinning site. In this case the vortex may be depinned and fall down into a neighboring pinning site. There is a critical generalized velocity \( u = u_c \) (\( u \) is proportional to the spin polarized current density) for which depinning occurs. In Figure 5.16 shows the displacement profile of the vortex core with clockwise chirality as a function of the AC frequency. The data is taken as the net displacement for a 100 ns AC pulse. The frequency range, for which propagation occurs, increases with increasing \( u \).

![Figure 5.16](image.png)

**Figure 5.16** Frequency dependent vortex core displacements for a 100 ns AC pulse. The nanowire thickness is 20 nm, the width is 240 nm, and the constriction wavelength is 480 nm. The constriction depth is 20 nm (a), 40 nm (b), 60 nm (c), and 80 nm (d). The vortex wall has a clockwise chirality.

For the specific wall type, the net displacement direction for \( d \geq 60 \text{ nm} \) is uniquely determined by its chirality and is in the positive \( x \) direction for a clockwise chirality and in the negative \( x \) direction for a counter clockwise chirality. For \( u > u_c \), the graphs show multiple peaks and the net displacement becomes more sensitive to the current frequency. This can be seen in Figure 5.17, where the apparent single peak in Figure 5.16 in the frequency interval 170 to 190 MHz in fact consists of three peaks. For the wires with \( d \leq 40 \text{ nm} \), the net displacement direction is dependent on both the chirality
and the frequency. The wire with $d = 20$ nm shows for $u = 80$ m/s and $f = 150$ MHz a displacement in the negative $x$ direction of $1.83 \mu m$. Increasing $u$ to 100 m/s gives zero or negative displacements for all frequencies. Two distinct peaks are present at 120 MHz and 150 MHz with respective displacements of -2.83 and -2.59 $\mu m$. In a wire with $d = 40$, a negative displacement is seen for $f = 140$ and 180 MHz, respectively. The net displacement direction is sensitive to the frequency, as can be seen when comparing the displacement of 2.63 $\mu m$ for $f = 170$ MHz with a displacement of -3.65 $\mu m$ for $f = 180$ MHz. This behavior can be understood if we treat the domain wall as a quasiparticle trapped in a potential well and consult Figure 5.9. For $d \geq 60$ nm, the amplification of resonant motion corresponds to the quasiparticle climbing up and down on alternating sides of the potential well, each time with an energy increase of $\Delta E_{dw}$, which corresponds to the energy absorbed by the domain wall from the driving current. If the potential heights of the barrier are $E_{b1}$ and $E_{b2}$ for negative and positive displacements, respectively, and $\Delta E_{b} = E_{b1} - E_{b2}$ the tendency for rightward displacement dominates if $\Delta E_{b} > \Delta E_{dw}$. The value of $\Delta E_{b}$ is 0.5, 1.1, 3.9, and 5.2 aJ for $d = 20$, 40, 60, and 80 nm, respectively.

![Figure 5.17](image)

**Figure 5.17** Net displacement of the vortex core as a function of AC frequency for $d = 80$ nm and $u = 100$ m/s. The pulse length of the AC signal is 100 ns.

### 5.7 Characterizing Domain Walls with MFM

All properties discussed here for 180° walls and in-plane domain magnetization, i.e. the type of wall (vortex or transverse), the axial magnetization alignment (head-to-head or tail-to-tail), the vortex chirality (clockwise or counter clockwise), and the vortex polarity (positive or negative), can be
determined by MFM. The experimental observations have been confirmed by numerical calculations. Each local magnetic moment in the simulated magnetic microstructure shown in Figure 5.18a produces a stray field according to expression (3.15). By integrating over the whole sample, a total stray field is obtained that is felt by a traversing MFM tip according to (3.14).

![Figure 5.18 A vortex wall in a nanowire. A magnetic vortex wall obtained from micromagnetic simulations (a). The expected MFM images calculated from the stray field generated by the magnetic structure in (a) for different scanning lift heights $h_{\text{lift}} = 10$ nm (b), $h_{\text{lift}} = 30$ nm (c), $h_{\text{lift}} = 50$ nm (d), $h_{\text{lift}} = 80$ nm (e), and a real MFM image with $h_{\text{lift}} = 50$ nm (f).](image)

In Figure 5.18b the MFM response is calculated for a constant tip-sample distance of 10 nm. In the centre of the vortex wall, a strong contrast is seen for a region with diameter of approximately 20 nm, corresponding to the vortex core with a magnetization pointing out from the plane. The region of spins aligned perpendicular to the wire axis, identified as the bright region in
Figure 5.18a, can be seen as a vague contrast forming an N-shaped pattern. The stray field from the edge roughness is seen as dots of alternating colors. Raising the lift height to 30 nm (Figure 5.18c) obscures some of the fine details. The vortex core and the edges seem smeared out and the N-shaped pattern is more pronounced. A head-to-head vortex wall with counter clockwise chirality can be identified with an N-shaped spin region. The given wall type with a clockwise chirality can be identified with a \( \Omega \)-shaped spin region (see paper I). For a lift height of 50 nm (Figure 5.18d), the vortex structure seems even more smeared out and the stray field from the edges is hardly visible. For a lift height of 80 nm (Figure 5.18e), the contrast from the vortex core is no longer discernable. In Figure 5.18f, an MFM image of a vortex captured at a lift height of 50 nm is shown. A head-to-head magnetization alignment of this vortex wall was controlled during the nucleation by applying a saturating field in the \(-y\) direction (see Figure 5.3) and then gradually reducing this field to zero. The chirality can be determined in Figure 5.18f from the N-shape and the small region of dark contrast in the centre of the vortex marks a vortex core with a positive polarity. The position of the observed vortex core coincides with the position of the simulated vortex core. The axial magnetization alignment can be determined from the contrast of the vortex domain wall. Reversing the axial magnetization alignment reverses the contrast of the domain wall. This reversal observed by MFM is indifferent to the chirality and was confirmed numerically.

5.8 Vortex Resonance

The dynamical properties of vortices in confined magnetic structures will be discussed in this section. A saturated magnetic structure responds only weakly to alternating driving forces. Magnetic structures with flux closure states, on the other hand, can be excited if the driving frequencies match the eigenfrequency of the system. Because of its gyrotropic nature, a magnetic vortex will enter a state of damped oscillations if perturbed from its equilibrium position. The translational mode can be calculated by Thiele’s equation of motion

\[
G \times \frac{dX}{dt} - \frac{\partial W}{\partial X} = 0, \quad (5.4)
\]

where \( G \) is the gyrovector, \( X \) is the position of the vortex core, and \( W = W(X) \) is the energy potential of the vortex [87]. The sense of gyration of the shifted vortex can be found from the gyrovector \( G = -2\pi pq e_z \) where \( p \) is the vortex polarity, \( q \) is the vorticity, and \( e_z \) is a unit vector in the \( z \) direction. Experimentally, the dependence of the vortex polarity on the sense of gyra-
tion has been confirmed by time resolved XMCD-PEEM measurements [101]. In this thesis, numerical calculations were performed on Permalloy nanowires and discs. Experimental measurements were carried out on a Permalloy disc with a diameter of 1 μm. The micromagnetic simulations were carried out using standard parameters for Permalloy: a damping parameter $\alpha = 0.02$, saturation magnetization $M_s = 860$ kA/m, and exchange stiffness $A = 13$ pJ/m. The cell size was chosen to be $5 \text{ nm} \times 5 \text{ nm} \times h \text{ nm}$.

![Image](image.png)

**Figure 5.19** Temporal evolution of the normalized magnetization $M_x$ of a vortex domain wall in a Permalloy nanowire (a). Amplitude of the FFT spectrum as a function of frequency for nanowires with different constriction depths $d$ (b).

Vortex domain walls trapped in nanowires with a thickness $h = 20$ nm, width $w = 240$ nm, constriction wavelength $l = 480$ nm, and constriction depth $d$ were investigated. The initial equilibrated micromagnetic state was obtained by subjecting the system to a static in-plane field of $\mu_0 H_A = 5$ mT in the $x$ direction. The system was then relaxed in zero field and the magnetization component $M_x$ recorded (see Figure 5.19a). The vortex core has a positive polarity and therefore orbits with a counter clockwise sense of gyration about the equilibrium position. By Fourier transforming the data in Figure 5.19a the corresponding resonance spectrum is obtained. The resonance spectra for $d = 20, 40, 60,$ and $80$ nm are shown in Figure 5.19b, which yield resonance frequencies of $285, 286, 325,$ and $303$ MHz, respectively.

The vortex dynamics in the discs were studied following essentially the same procedure as for the nanowires. The eigenfrequencies determined from resonance spectra of discs with different geometry are plotted in Figure 5.20. For a given thickness, the eigenfrequency decreases with increasing diameter. For a given diameter, the eigenfrequency increases with increasing thickness. The experimental value of the eigenfrequency was found to be $284$ MHz for a disc with $(40 \pm 10)$ nm thickness and $1$ μm diameter, measured using the homodyne detection method (see section 3.3.3). The simulated eigenfrequencies for a disc with $1$ μm diameter and $h = 30, 40,$ and $50$ nm are $268, 344,$ and $416$ MHz, respectively.
Figure 5.20 Resonance frequency in Permalloy discs as a function of the disc diameter. The data are shown for simulations (filled circles) of discs with thickness $h = 30, 40, \text{ and } 50 \text{ nm}$. 
The class of materials where magnetism is incorporated by a fractional substitution of atoms in a semiconducting host is called diluted magnetic semiconductors (DMS). These materials are especially interesting since the coexistence of magnetism and semiconductor properties could enable one extra degree of freedom in future electronic devices. The subclasses of DMS commonly considered are based on the II-VI, III-V, IV-VI, or IV systems where the Roman numerals notify which groups in the periodic table the constituent semiconductor elements belong to.

Semiconductors are characterized by their sensitivity to impurity atoms that changes the electrical conductivity. A doping material can therefore be used to modify the number of free charge carriers in the semiconductor. Depending on the valence of the dopant, the doped semiconductor is either of n-type if the number of negative charge carriers (electrons) is increased or of p-type if the number of positive charge carriers (holes) is increased.

In DMS, a transition metal, e.g. Ni, Co, Fe, Gd, or Mn, with an atomic concentration $x$ is used as the magnetically contributing dopant. This dopant can embody the semiconductor in two different ways, either by substituting one of the host cations or by occupying an interstitial position in the lattice. Mn has proved to be advantageous in several aspects. First, it may be introduced into a II-VI host in high fractions, up to 50% in selenides and up to 80% in tellurides [102], without significantly altering the crystal quality. In III-V hosts, this fraction is lower, typically below 10%. However, successful growth of In$_{1-x}$Mn$_x$As with $x = 0.18$ has been reported [103]. Secondly, the localized 3d-electrons of the Mn$^{2+}$ ion provide a high spin state of $S = 5/2$. This corresponds to a magnetic moment of between 4 and 5 $\mu_B$ per uncompensated substitutional Mn (Mn$_{Ga}$) in Ga$_{1-x}$Mn$_x$As for fractions between $x = 0.017$ and $x = 0.067$ [104]. Furthermore, in II-VI hosts the valence of Mn matches that of the cation so it becomes neither an acceptor nor a donor.

The II-VI based compounds, such as Cd$_{1-x}$Mn$_x$Te and Zn$_{1-x}$Mn$_x$Se, have been studied in detail over the last decades [105]. They are easy to fabricate both in bulk form and as thin films. However, because of doping difficulties many of these materials have lost popularity from an application point of view [106]. Instead, the III-V based semiconductors have become increasingly popular. This is mainly due to the refined fabrication techniques that came with low temperature molecular beam epitaxy (LT-MBE) and allowed orders of magnitude larger doping concentrations to be incorporated [107].
should be mentioned that other subclasses of DMS exist. Transition-metal doped IV-VI compounds such as Ge$_{1-x}$Mn$_x$Te [108], Pb$_{1-x}$Mn$_x$Te, and Pb$_{1-x}$Gd$_x$Se [109] have been studied. However, the exchange interaction constant for these materials is small, typically one order of magnitude smaller as compared to II-VI and III-V compounds [110, 111]. Recently, increasing attention has been directed towards other compatible candidates among the doped IV systems, for example MnGe [112, 113] or CoFeSi [114].

In this work, the III-V system Ga$_{1-x}$Mn$_x$As has been studied. This system belongs to one of the most well understood systems in terms of combined magnetic and semiconductor properties and therefore serves as a suitable DMS prototype material. Similar to several other III-V systems, Ga$_{1-x}$Mn$_x$As crystallizes in a zinc-blende structure shown in Figure 6.1. The lattice parameter $a(x)$ depends on the dopant concentration $x$ according to $a(x) = 0.566 + 0.032x$ nm [115]. Here, the films were grown on GaAs (001) substrates and the magnetic properties were probed with a SQUID MPMSXL magnetometer. This device is necessary for measuring the tiny magnetic signals from thin films with a low doping concentration of magnetic elements. In fact, the ferromagnetic response from the sample is at high magnetic fields often inferior to the magnetic contribution from the substrate.

### 6.1 The Ferromagnetic Transition Temperature

In order for DMS to be useful in technological devices the critical temperature $T_c$ at which the transition from paramagnetism to ferromagnetism occurs needs to be higher than room temperature. During the last two decades, $T_c$ has been substantially increased by scientific progress. A key factor in this progress is that the Mn content in III-V based InMnAs and GaMnAs alloys has a large impact on $T_c$. In 1992 the first successful LT-MBE growth of InMnAs with an Mn content above 1% yielded a $T_c$ of 7.5 K [107]. In 1997 the first ferromagnetic alloys of GaMnAs were reported with a $T_c$ of about 50 K [116]. By optimizing the growth conditions and Mn content the record was set to 110 K in 1998 [106]. The next jump in $T_c$ to 173 K as reported in 2005 was to a high degree accredited to the application of post growth annealing techniques [117]. In the IV-based CoFeSi, a transition temperature of 53 K has been reported [114].
Figure 6.1 The crystal structure of the semiconductor GaAs. If Mn is introduced as a dopant, three different scenarios are possible: (1) it can substitute the Ga atom; (2) it can find an interstitial position between four cations; or (3) it can find an interstitial position between four anions.

6.2 Defects in Ga$_{1-x}$Mn$_x$As

Each substitutional Mn$_{Ga}$ in Ga$_{1-x}$Mn$_x$As acts as a monovalent acceptor and provides a local magnetic moment with zero angular momentum and spin $S = 5/2$. Nevertheless, several defects are commonly present that influence the electrical and magnetic properties of the system. A Ga can for example be replaced by As. This As antisite As$_{Ga}$ acts as a double donor and contributes to hole compensation. The defect originates from the high As pressure, which is used at LT-MBE to assure the layer-by-layer growth without Ga droplets forming on the surface. Due to the high thermal stability of As$_{Ga}$, up to $\sim 450 \, ^\circ$C [118], post growth annealing is not a viable option since Mn tends to form secondary phases with As even at much lower temperatures. One way to reduce the concentration of antisites is to use As$_2$ dimers instead of As$_4$ tetramers during film growth [119].

Another defect is the Mn interstitial Mn$_I$ that also acts as a double donor. Because of the attractive Coulomb interaction between negatively charged Mn$_{Ga}$ and positively charged Mn$_I$ the interstitials are expected to form pairs with the substitutional Mn during the growth [120]. It has been suggested, both from experiments [121] and from theoretical calculations [122], that the Mn$_{Ga}$-Mn$_I$ pair couples antiferromagnetically and therefore provides a com-
pensated local spin, which in experiments is reflected in a decrease of the measured saturation magnetization. Mn$_1$ can occupy a position surrounded by four anions, or a position surrounded by four cations. These two positions result in different local energy landscapes and it could therefore be argued that the interstitial has a spatial preference for one of the situations. Nevertheless, the situation is not yet completely understood as some report Mn$_1$ to favor a position between cations [120, 122], whereas other conclude that Mn$_1$ favors a position between anions [123-127].

![Graph](image)

**Figure 6.2** The influence of post growth annealing on the magnetic properties. The critical temperature $T_c$ is shown in (a) as a function of the annealing time $t_a$ for the annealing temperatures $T_a = 190$, 215, and 240 °C. As seen in the inset the saturation magnetic moment increases with increasing $t_a$. The magnetic moment, measured in Bohr magnetons $\mu_B$ per Mn-ion, as a function of temperature is shown for $T_a = 215$ °C in (b) and $T_a = 240$ °C in (c). The symbols connected to the graphs in (b) and (c) correspond to different annealing times.

### 6.3 Enhancement of $T_c$ by Post Growth Annealing

Annealing Ga$_{1-x}$Mn$_x$As films at high temperatures leads to the formation of MnAs clusters [116]. However, annealing at temperatures about the growth temperature or slightly below can substantially elevate the critical temperature of the material [128-130] due to the diffusion of Mn$_1$ to the free surface.
followed by its passivation with oxygen or nitrogen. In paper VI, the influence of the annealing parameters on the magnetic properties is investigated. The involved parameters that are varied in this study are the annealing temperature $T_a$ and the annealing time $t_a$. In principle, the growth temperature $T_g$ and the Mn concentration $x$ could also be varied. However, these parameters are kept constant at $T_g = 230 \, ^\circ\text{C}$ and $x = 0.06$ for all samples. As seen in Figure 6.2a, $T_c$ has a clear dependence on annealing time. For $t_a > 13$ hours, $T_c$ enters a region where it changes only slightly with increasing time. This tendency is nearly the same for the different annealing temperatures and infers that an increase of $T_c$ of approximately 65% compared to the as-grown films could easily be obtained by choosing a sufficiently high $t_a$. For $0 < t_a < 13$ hours, on the other hand, a different behavior is shown that more depends on $T_a$. By optimizing the parameters $t_a$ and $T_a$, an increase in $T_c$ of about 85% is possible to achieve. The inset shows the saturation magnetization dependence on $t_a$ and $T_a$. Thus, by assuming that MnI has a compensating effect on the magnetic moment the inset hints at the passivation of MnI by annealing time. A layer of amorphous As is used as a capping layer for the epitaxially grown Ga$_{1-x}$Mn$_x$As films. The main role of this capping layer is to act as a reactive medium for the passivation of MnI. Experimentally, the diffusion of Mn was confirmed by secondary ion mass spectroscopy (SIMS). A Mn intensity profile, obtained by SIMS, is shown in Figure 6.3. As expected, the passivated MnI in the capping layer can be seen in the annealed samples but not in the as-grown samples.

![Mn intensity profiles](image.png)

**Figure 6.3** Mn intensity profiles of a Ga$_{1-x}$Mn$_x$As film obtained by SIMS for the annealing times $t_a = 0.5$, 2, and 20 hours. The film is 1000 Å thick and capped with an amorphous As layer. The annealing temperature is 240 °C. The inset shows a close up of the intensities at smaller sputter depths. Notice that the As cap desorbs for annealing times between 0.5 h and 2 h.
6.4 Origin of Ferromagnetism in Ga$_{1-x}$Mn$_x$As

The origin ferromagnetism in Ga$_{1-x}$Mn$_x$As can be explained using contributions from several interactions [104, 115, 133]. The net effect of these contributions is to a large degree dependant on the dopant concentration $x$. For $x < 0.01$ ferromagnetism has not been experimentally observed. At these concentrations, the average distance $r_c$ between Mn impurities is larger than the extent of each hole bound to Mn, which can be approximated by the Bohr radius, $r^* = \varepsilon h/(2\pi m^* e^2)$, where $\varepsilon$ is the dielectric constant of the material and $m^*$ is the effective hole mass. Ferromagnetism has been measured for $x$ about 0.01 [115]. At these concentrations, $r_c$ and $r^*$ are of the same order and holes can mediate ferromagnetism between the Mn moments. This mechanism can be described by Zener’s double exchange interaction [134], which employs an intermediate nonmagnetic atom through which two separated Mn atoms in the impurity band can exchange 3d-electrons. A requirement here is that the two atoms have a different number of electrons in their respective 3d-shell. From Hund’s rule, parallel spin alignment is favored in the 3d-shell. This decreases the energy of the spin polarized electrons and increases the probability of electron hopping [104]. For higher concentrations the impurity band merges with the valence band. The coupling between local Mn moments and carrier spins can be described by the Ruderman-Kittel-Kasuya-Yosida (RKKY) interaction [111]. An itinerant carrier is subjected to exchange interaction with a local moment and becomes spin polarized. Upon scattering with another Mn site, this carrier couples its spin to the local moment. Depending on the distance $r$ between the local Mn moments the interaction is either of ferromagnetic or antiferromagnetic character. More specifically, the interaction between two local moments $S_1$ and $S_2$ is described by the Heisenberg-like Hamiltonian,

$$H_{RKKY} = -J_0 F(2k_F r) S_1 \cdot S_2,$$

(6.1)

where $k_F$ is the Fermi wave vector and $J_0$ is an interaction parameter that depends on the effective mass of the carrier. The function $F(x)$ oscillates according to

$$F(x) = \frac{1}{x^4} (x \cos x - \sin x).$$

(6.2)

The oscillatory nature of the RKKY Hamiltonian thus allows a coexistence of ferromagnetic and antiferromagnetic interactions in the system.
7 Controlled Motion of Magnetic Particles

In this section we will discuss how patterned elliptical Permalloy microstructures can be used in an application for controlled motion of magnetic particles. The main motivation of this study comes from the existing techniques that deal with the magnetic separation of biomolecules in a suspension. Magnetic separation is an already widely employed technique where magnetic particles with functionalized surfaces bind to target molecules [135-138]. The magnetic particles together with the bound targets can then be separated from the unbound biomolecules using external magnetic fields. Since this method makes no distinction between the magnetic particles, they are all separated as a group from the suspension. Manipulation of individual magnetic particles would make possible the use of more sophisticated particle sorting mechanisms.

Trapping and untrapping of magnetic particles on patterned magnetic elements can be achieved by using external magnetic fields and the stray fields generated by the magnetized elements [139]. Transport of magnetic particles has been demonstrated using the stray field gradients from electrically conducting wires [140]. In papers VII-IX, a technique is discussed that allows us to trap magnetic particles, transport them along patterned arrays, and selectively separate individual particles [141].

Figure 7.1 A SEM image of a magnetic bead sitting on a magnetized elliptical Permalloy element. The diameter of the bead is 2.8 μm and the lateral size of the element, which is not fully shown, is 2 μm × 6 μm.
7.1 The Magnetic Particles

The magnetic particles, or beads, used here are spheres made of porous polystyrene with embedded iron oxide nanoparticles [142]. The beads are monodisperse and have a diameter of 2.8 μm. In an applied magnetic field these beads have a superparamagnetic response and the reorientation of the magnetic moment with the applied field is within the embedded iron oxide nanoparticles. A magnetized magnetic element will attract the bead with its stray field. Figure 7.1 shows a SEM image of a bead immobilized at the rim of an elliptical element. The beads are coated with functionalized polyethylene glycol (PEG-NH₂) with amino groups that are responsible for the conjugation to target proteins. A suspension of the beads, based on ionized water, is then prepared. Due to the negligible bead remanence, magnetic attraction between the beads in zero external magnetic field is not a major problem. However, some conglomeration may occur due to non-magnetic forces (e.g. electrostatic interactions, hydrophobic interactions, and hydrogen bonds) between the beads. This force can in some cases be very large and dominate the magnetic force by orders of magnitude [143]. Even without functionalized surfaces some aggregation of the beads may occur. To amend the situation a small quantity of detergent is added to the suspension.

Figure 7.2 A part of the transport line for 2 μm × 6 μm elements where adjoining ellipses have their long axes along orthogonal directions. The separation between two adjoining elements is 3 μm. In (a) the defined transport line is shown. In (b) an MFM image of the fabricated elements is shown where a magnetic field of 4kA/m is applied in the negative x direction. Elliptical elements having their easy axis parallel to the field are saturated, indicated by the dark and bright ends. Elements with their hard axis parallel to the field exhibit multidomain states.
7.2 The Elliptical Elements

Elliptical Permalloy elements with a thickness of 50 nm and two different lateral geometries, 1 μm × 3 μm and 2 μm × 6 μm, were fabricated using electron beam lithography, material deposition by thermal evaporation, and a final chemical lift-off step. The element sizes were chosen to be small enough for handling the type of magnetic particles commonly used in magnetic separation techniques [144] and large enough to allow the magnetic microstructure to be in a multidomain state. The aspect ratio of 1:3 was chosen to provide multidomain states for certain applied field directions and saturated single domain states for other field directions. The field necessary to saturate the element along different directions depends on the aspect ratio. Along the short axis of the ellipse, the field needed for saturation increases with increasing aspect ratio. For saturation along the long axis, the saturation field decreases with increasing aspect ratio. A magnetized elliptical element will attract a bead in the vicinity towards one of its poles. The size of the attractive force can be estimated by considering an element magnetized along an easy axis and a bead free to move parallel to this axis. The calculated force is 0.2 nN at a separation of 3 μm and about 3 nN at the contact point [145]. This attractive force between an element and a bead increases with the net magnetization of the element.

The elements are arranged in arrays, or transport lines, according to Figure 7.2a. Each element is separated from its orthogonal neighbor by a distance of 3 μm. In Figure 7.2b an external field is applied in the -x direction. The horizontal elements, having the long axes along the x direction, are saturated and generate strong stray fields that can attract a bead in the vicinity. The vertical elements, with long axes along the y direction show multidomain states. These states are of flux closure type and generate weak stray fields. Nevertheless, they have a small net magnetization in the direction of the applied field.

7.3 Transport and Separation Principles

The motion of a bead is initialized and maintained by an applied magnetic field that is rotating in plane with constant magnitude. Here, a requirement for the operation is that each element can alternate between a multidomain state and a single domain state as the external magnetic field is rotated. Under this rotation, a trapped bead will follow the perimeter of an element according to Figure 7.3. As the bead gets into the region of shortest distance between two elements, the element on which the bead is sitting will be in a flux closure state. At the same time, the neighboring element will be in a saturated state and will thus generate a strong stray field that attracts the bead. The bead will now jump to the neighboring element and continue its
motion. By changing the sense of rotation of the field, the direction of motion of a bead about the element is correspondingly reversed. However, the net displacement and direction of a bead along a transport line remains the same regardless of the sense of field rotation.

**Figure 7.3** The motion of a bead as the applied field is rotated counter clockwise. The white and black circles indicate the starting and end positions of the bead, respectively. The bead follows the direction of the net magnetization about the element’s circumference. As the bead reaches the region of shortest distance between an element and its neighbor, the former will be in a flux closure state and the stray field from the latter will attract the bead. The bead jumps to the neighboring element and continues its motion along the dashed path.

**Figure 7.4** Photographs showing a junction formed by adding two orthogonal transport lines of 2 μm × 6 μm elements. Two cases are considered where the choice of path at the junction is determined by the sense of rotation of the applied field. The hollow and solid black circles illustrate the start and final positions of the bead, respectively. (a) The field is rotated clockwise and the bead propagates to the left at the junction. (b) The field is rotated counter clockwise and the bead propagates straightforward.
A junction is formed by connecting two transport lines orthogonally as shown in Figure 7.4. Here, the sense of rotation of the field determines the choice of path at the junction. A clockwise rotating field will make the bead turn left at the junction, whereas a bead will continue straightforward for an anticlockwise rotating field. Single particles can now be selectively separated by temporarily reversing the sense of field rotation.

**Figure 7.5** Micromagnetic simulation of the domain configuration in a 1 μm × 3 μm sized element as the applied field is rotated. The applied field strength is $\mu_0 H_A = 15$ mT. (I) Rotating the field from 0° to 60° only shifts the net magnetization angle by 13°. (II) A multidomain state appears at $\theta_e = 90°$ for which the magnetization is 0.44$M_s$. (III) At $\theta_e = 100°$ the multidomain state is still present and the magnetization angle precedes $\theta_e$. (IV) The domain configuration is nearly saturated with a net magnetization of 0.97$M_s$. 

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The detailed evolution of the magnetization in an element as the applied field, with magnitude $H_A$, is rotated can be studied using micromagnetic simulations. For an element being in a perfectly saturated magnetic state, the single domain state transforms to a multidomain state as the applied field direction $\theta_e$ reaches a certain nucleation angle $\theta_n$. At this point, the net magnetization decreases abruptly. For fields not sufficiently large to saturate the structure along $\theta_e = 0$ the magnetic configuration rather seems to be in a nearly saturated state with a vortex structure still present at each end of the long axis. Therefore, upon rotation of the applied field there is no abrupt transformation but instead a gradual propagation of the vortices into the central part of the element. In such a case, the nucleation angle can be approximated as the angle for which the net magnetization decreases below $0.5M_s$. The nucleation angle increases with increasing $H_A$. After nucleation the multidomain structure remains for an angular range, which decreases with increasing field $H_A$. For the structure in Figure 7.5 using $\mu_0H_A = 15$ mT this angular range is approximately $50^\circ$ wide. Generally, the net magnetization direction lags $\theta_e$ with an angle $\phi = \phi(\theta_e,H_A)$ for $0 < \theta_e < \theta_n$ and precedes $\theta_e$ for $\theta_n < \theta_e < 180^\circ$.

**Figure 7.6** The experimental setup. (A) An optical microscope. (B) Fluid inlet and outlet for the suspension of magnetic beads. (C) Electrical connection to the electromagnet. (D) Fluid cell. (E) Sample with patterned magnetic arrays mounted inside the fluid cell.
In Figure 7.6, the experimental setup is shown. The sample substrate is positioned in the centre of a fluid cell. This fluid cell has a transparent top side and two drilled holes for the inlet and outlet of the suspension. A vector electromagnet (not visible in the figure) is located under the fluid cell. The magnet can generate fields up to approximately 40 kA/m with arbitrary in-plane directions. Typically, the field is gradually rotated with angular steps of a few degrees and constant amplitude of about 6 kA/m for the 2 µm x 6 µm elements. An optical microscope, placed just above the fluid cell, is used to monitor the motion of the particles.
Summary in Swedish

Inledning
Magnetiska tunna filmer har andra egenskaper än samma material i bulk-form. En films magnetiska egenskaper är starkt beroende av dess storlek och form vilket kan utnyttjas när man skräddarsyr de magnetiska egenskaperna för specifika ändamål. Detta har stor betydelse inom datalagringsindustrin och den snabba utveckling som pågår är direkt kopplad till de framsteg som gjorts inom forskning på magnetiska tunna filmer.

Fysikalisk förståelse för ett material kan erhållas genom att tillverka materialet och empiriskt studera dess egenskaper. Om den fysikaliska modellen redan finns kan man utföra mikromagnetiska beräkningar på systemet för att bestämma dess magnetiska jämviktstillstånd och härigenom studera hur materialets magnetism delar upp sig i små områden (magnetiska domäner) samt hur systemet kan manipuleras med yttre krafter. Materialets geometri och magnetiska egenskaper kan därför bestämmas innan man beslutar sig för att deponera och mönstra filmerna. Domänstrukturen i de färdiga filmerna kan undersökas med ett flertal metoder, t ex med hjälp av MFM\(^1\) eller XMCD-PEEM\(^2\), och man finner då att de experimentella resultaten i hög grad stämmer överens med de numeriska beräkningarna. En annan fördel med de numeriska beräkningarna är att man kan få fullständig information om systemets spinnkonfiguration för varje tidssteg, som typiskt är av storleksordningen pikosekunder (en tusendel av en miljarddels sekund). Detta medför att man kan studera tidsberoende magnetiska egenskaper hos ett material.

Domänväggar
I en tunn film, med en tjocklek på några tiotals nanometer, av ett mjukmagnetiskt material så föredrar magnetiseringen att ligga i planet. Reducerar man dimensionaliteten ytterligare för detta material erhålls en tråd vars magnetisering strävar efter att ligga parallellt med tråden. Gränsskiktet där två domäner med motriktade magnetiseringar möts utgörs av en domänvägg. Åtskillnad görs generellt mellan två typer av magnetiska domänväggar. I en transversell domänvägg pekar magnetiseringen i planet vinkelrätt mot trådens axel. I en vortexvägg däremot cirkulerar magnetiseringen runt ett om-

\(^1\) Magnetic force microscopy.
\(^2\) X-ray magnetic circular dichroism photo electron emission microscopy.
råde, vortexkärnan, där magnetiseringen är riktad vinkelrätt mot planet (se Figur 1). Vortexdomänväggar har studerats i ett flertal arbeten [146-149], speciellt med avseende på hur vortexkännans magnetiseringsriktning påverkar dess dynamik. Dock saknas full förståelse om hur vortexens kiralitet, som är positiv (negativ) för medsols (motsols) magnetiseringsfördelning runt vortexkärnan, påverkar de magnetiska egenskaperna. I denna avhandling studeras hur och var domänväggar fastnar i platta nanotrådar med artificiellt beskurna kanter (sågtandsnanotrådar) och hur dessa så kallade pinningegenskaper beror på vortexens kiralitet. Med hjälp av mikromagnetiska beräkningar har energipotentialen för en vortexvägg i ett pinningcentrum beräknats för de olika kiraliteterna. Resultaten visar att energipotentialen i en given nanotråd uppvisar en asymmetri som tilltar med djupet på beskärningen. Som visas i artikel I kan energibarriären höjd direkt bestämmas med ett pålagt magnetfält.

Spinnrörelsemängdsmomentet hos en spinnpolariserad ström kan överföras till de lokala magnetiska momenten i domänväggen som därigenom förflyttas i elektronernas rörelseriktning. En spinnpolariserad likström kan förskjuta domänväggar med reduce rad rörelseriktning. Experiment som utförts med likström har visat att spinnpolariserad växelström kan förflytta domänväggar med reducerad strömtäthet [100]. Om växelström appliceras i en rak tråd kommer domänväggens nettoförflyttning att vara noll eftersom elektronernas rörelseriktning alternarar. Om växelström däremot appliceras i en sågtandsnanotråd kan domänväggens förflyttning med en likriktad nettoförskjutning vars riktning bestäms av kiralitetens typ. Detta är möjligt då den resonansmässigt förstärkta rörelsen överskrids den lägre barriären. Domänväggars egenskaper i magnetiska nanotrådar studeras i artiklarna I-III.

Anisotropi

För att förstå ett magnetiskt material måste man studera dess magnetiska anisotropi. I ett material som är isotropt finns inget riktningssberoende, d.v.s. de magnetiska egenskaperna ser likadana ut åt alla håll. Detta klassas dock som ett undantag eller en idealiserad modell snarare än ett realistiskt exempel. Magnetisk anisotropi yttrar sig i olika former och man skiljer normalt på den magnetokristallina anisotropin, som hänrrör till materialets kristallstruktur, och den inducerade anisotropin som kommer av avvikelser från den ideala kristallstrukturen. Formanisotropi är en annan term som kopplar till materialets geometri och gör det att magnetiseringen i mjukmagnetiska tunna filmr lägger sig i planet. I vissa material dominerar ut-ur-plan anisotropin, även i tunna filmr. Dessa är särskilt intressanta för applikationer inom datalagring eftersom de tillåter en hög informationstäthet. Legeringar av FeCo med stark ut-ur-plan anisotropi studeras i artiklarna IV och V.
Spinntronik och utspädda magnetiska halvledare


Fältinducerad manipulation av mikrometerstora magnetiska partiklar

Magnetisk separation av biomolekyler i heterogena lösningar är en väletablerrad teknik inom bioteknologin. Magnetiska pärlor med funktionaliserade ytor binder selektivt till en målgrupp av biomolekyler och kan separeras från lösningen med hjälp av pålagda magnetfält. Således separeras pärlor och biomolekyler som en enhet från lösningen. En metod där biofunktionaliserade pärlor kan separeras individuellt skulle göra det möjligt att behandla extremt små mängder av biomolekyler vilket i sin tur kan generera nya tillämpningar inom så kallad lab-on-a-chip teknologi. I denna avhandling diskuteras en metod för att förflytta superparamagnetiska partiklar längs mönstrade elliptiska magnetiska mikrostrukturer med hjälp av pålagda roterande magnetfält. Man utnyttjar dock inte det pålagda fältet direkt, utan istället det strömfält som genereras från de elliptiska mjukmagnetiska elementen av

³ Magnetoresistive random access memory.
⁴ Spin transfer torque MRAM.
Fe$_{20}$Ni$_{80}$ (permalloy) i ett pålagt homogent magnetfält. De magnetiska elementen placeras i rad i ett särskilt mönster så att de i ett roterande magnetfält magnetiseras på ett sådant sätt att de magnetiska pärlorna kontrollerat förflyttas längs raden av element. Separation av pärlor uppnås genom att ortogonalt koppla samman rader av mönstrade elliptiska element. Experimentella studier och mikromagnetiska simuleringar av domänstrukturen för de elliptiska elementen som visar för vilka parametrar partikeltransporten kan ske beskrivs i artiklarna VII-IX.

**Figur 1** Vortexkärnans ut-ur-planet magnetisering i en 240 nm bred nanotråd erhållen från numeriska beräkningar. Systemet är uppdela i ett nät av beräkningsceller (syns i figuren) med cellstorlek 5 nm $\times$ 5 nm. I figurens övre vänstra hörn syns i-plane-magnetiseringskonfigurationen (pilar) för en vortexdomänvägg, med motsols kiralitet, i en nanotråd.
First, I would like to thank my supervisors Peter Svedlindh and Per Nordblad for all support, guidance, knowledge, and always having time. Indeed, the research here at the division of Solid State Physics has been both exciting and enjoyable.

I am thankful to all members of the magnetism group (and closely related people) for nice collaboration and companionship, Klas Gunnarsson for introducing me to the MFM and being a cheerful office neighbor (and towing my car a cold and rainy day), Magnus Wikberg for exploring Madrid together during the Intermag conference, Rebecca Bejhed for appreciated initiative and organizational skills, Matthias Hudl for being one of the driving forces behind the magnetism group’s get-togethers, Hans Lidbaum for discussions and producing SEM images for the biomagnetic project, Roland Mathieu and Petra Jönsson for discussions and advice concerning research in Japan, Gabriella Andersson for cooperation and discussions on out-of-plane materials, Andreas Liebig for teaching me how to deposit films using “Binford”. All members of the Friday-article group are acknowledged for interesting discussions on various magnetic topics.

I would also like to thank former members of the magnetism group who were a source of inspiration when I started my PhD studies: Victor Stanciu for introduction to the cryogenic world and all culinary gatherings, Solveig Felton for enjoyable discussions, and Pierre Roy for undamped discussions on micromagnetism.

In our modern society, unexpected computer failure is not a question of “if”, but “when” it will happen. In these times of despair Jonatan Bagge has been very helpful in restoring or improving the situation. I also thank Mikael Österberg for setting up and maintaining remote desktop connections and Enrique Carrasco for providing the latest software to my computer.

I express my gratitude to Claes-Göran Granqvist for valuable advice of financial character, Bengt Götesson for engineering skills and advice concerning computer graphics, Inger Ekberg for helping me with administrative challenges, Ingrid Ringård for always making me feel welcome to discuss things, and the rest of the administrative group at the department for being kind and helpful.

I would like to express my gratitude to Christian Sondell for producing graphic material for some presentations and Liliana Arevalo for expertise on FEM calculations.
Part of this work was carried out at the Institute for Chemical Research at Kyoto University. I am grateful to Teruo Ono for generously inviting me to the nanospintronics lab (and of course all the tennis games!). I would like to thank all members of the lab (including the technician and the new students) for not only taking time to explain the experimental equipment but also eating out, playing sports, having fun, and creating a wonderful atmosphere: Kensuke Kobayashi for helping me to obtain civil rights and settle down upon arrival, Shinya Kasai for being a walking library of information, Atsushi Himeno for taking us far far away by car just to find the best ramen, Hironobu Tanigawa for the wedding party, Masayuki Hashisaka for the nice jokes, Kuniaki Yano for unforgettable entertainment, Michael Delmo for helpful translations, Kouta Kondou for never saying no, Keisuke Yamada for initiating sport events, Yoshinori Tamada for being neighbors in the office, Shuji Nakamura for the new year card, Takanori Inoue for challenging games of table tennis, and our secretary Hiroko Toyama for maintaining order in the lab and creating a joyful atmosphere. In connection to this I would also like to thank Yoshinobu Nakatani from Tokyo for the micromagnetic discussions and simulation data.

I am indebted to Yasusei Yamada for all help before arriving to Japan and hospitality at Mikawa Anjo (and for planting rice together).

Part of this work was performed in the nanomagnetism research group at the University of Konstanz. I gratefully acknowledge Mathias Kläui for inviting me to the lab and the events (scientific and social) taking place there. My gratitude goes to all members of the lab for providing a friendly and productive atmosphere, Ulrich Rüdiger for interesting discussions, Olivier Boulle for collaboration on the spin-torque measurements, June-Seo Kim for sharing many hours in the lab (and investigating applied forms of spin transfer), Lutz Heyne and Jan Rhensius for the night shifts at BESSY, Stephen Krzyk for helping me with the SEM, Jan Nievendick for interesting discussions and a nice trip to the Alps, Jan Heinen for the spider, and Christian Dette for cooperative efforts on the cryogenic sample rod and exploring the city together.

The time here at the Solid State Physics division would not have been the same without all nice colleagues (present and former) sharing lunches, coffee breaks, TMCPs, Christmas dinners, barbeques, and social events. Thank you all for creating a pleasant atmosphere! I would also like to thank Barbara Piechocinska for waking up my mind, Magdalena Lundh for the avocado plant contest, Tobias Boström for a nice time in Freiburg and on the climbing wall, and Zareh Topalian for enjoyable discussions (and games of chess).

Finally, I would to thank my family and friends for always being there.
Bibliography


A doctoral dissertation from the Faculty of Science and Technology, Uppsala University, is usually a summary of a number of papers. A few copies of the complete dissertation are kept at major Swedish research libraries, while the summary alone is distributed internationally through the series Digital Comprehensive Summaries of Uppsala Dissertations from the Faculty of Science and Technology. (Prior to January, 2005, the series was published under the title “Comprehensive Summaries of Uppsala Dissertations from the Faculty of Science and Technology”.)