Spin transfer torques and spin dynamics in point contacts and spin-flop tunnel junctions

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Abstract

The first part of this thesis is an experimental study of the spin-dependent transport in magnetic point contacts. Nano-contacts are produced micromechanically, by bringing a sharpened non-magnetic (N) tip into contact with a ferromagnetic (F) film. The magnetic and magneto-transport properties of such N/F nanocontacts are studied using transport spectroscopy, spanning the ballistic, diffusive, and thermal transport regimes.

Single N/F interfaces can exhibit current driven magnetic excitations, which are often manifest as peaks in the differential resistance of a point contact defining the N/F interface. Our experiments show that such surface magnetization excitations, and thus the single-interface spin torques, are observed for diffusive and thermal transport regimes where the conduction electrons experience strong scattering near the N/F interface, and are absent for purely ballistic contacts. We conclude that the single-interface spin torque effect is due to impurity scattering at N/F interfaces.

Single N/F interfaces can also exhibit hysteretic conductivity, which is qualitatively similar to the spin-valve effect found in F/N/F trilayers. Based on our measurements of N/F point contacts in the size range of 1-30 nm, we propose two mechanisms of the observed hysteresis. The first mechanism relies on a non-uniform spin distribution near the contact core and is magnetoelastic in origin. This interpretation is in good agreement with some of our experiments on larger point contacts as well as with a numerical micromagnetic model we have developed, where a stress-induced anisotropy creates a non-uniform, domain-wall-like spin distribution in the contact core. The second mechanism we propose is a surface effect which relies on a difference between the surface and interior spins in the ferromagnet in terms of their exchange and anisotropy properties. The surface spin-valve mechanism is in good agreement with the hysteretic magnetoresistance observed for our smallest contacts (∼ 1 nm) and for contacts to nanometer thin ferromagnetic films. This interpretation means that the surface magnetization can be reduced and weakly coupled to the interior spins in the ferromagnet. We find that this surface spin layer can be affected by both external fields and the spin torque of a transport current. The surface magnetization can even form nano-sized spin vortices at the interface.

The nature of the magnetic excitations induced by by nominally unpolarized currents through single N/F interfaces was probed directly using microwave irradiation. We observed two characteristic high-frequency effects: a resonant stimulation of spin-wave modes by microwaves, and a rectification of off-resonant microwave currents by spin-wave nonlinearities in the point contact conductance. These experiments demonstrate that the effects observed are spin-dynamic in nature.

In the second part of the thesis we study the spin-dynamics in spin-flop tunnel junctions used in toggle magnetic random access memory. Current pulses in the range of 100 ps used to excite the magnetic moments of the two coupled Py free layers into an oscillatory state, in both the antiparallel and scissor states of the cell. These oscillations are detected directly by measuring the junction resistance in real time with a 6 GHz measurement bandwidth. The junctions had the shape of an ellipse, with lateral size ranging from 350 × 420 to 400 × 560 nm. The optical and acoustical precession modes of the the spin-flop trilayer are observed in experiment, as expected from single-domain model. The experimental spectra contain additional features, which are explained using numerical micromagnetic simulations, as originating from magnetic state transitions between different magnetization states with non-uniform spin distributions.
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TO MY FAMILY
List of abbreviations

AC  alternating current
AF  antiferromagnetic
AMR  anisotropic magnetoresistance
AP  antiparallel state or alignment of the spins, magnetization etc.
BL  bit line
CDHS  current-driven hysteretic switching
ch.  chapter in the manuscript
CIMR, (CIMS or CIS)  current-induced magnetization reversal (switching)
CIP  current in plain
CPP  current perpendicular to the plain
DC  direct current
DMR  domain wall magnetoresistance
DOS  density of states
DW  domain wall
EA  easy axis
F, FM  ferromagnet, ferromagnetic
FDHS  magnetic field-driven hysteretic switching
HA  hard axis
HF  high frequency
IVC  current-voltage characteristics
LA  longitudinal
µMS  micromagnetic simulations
MR  magnetoresistance
MRAM  magnetic random access memory
MTJ  magnetic tunnel junction
N  non-magnetic
ND  nano-domain
P  parallel state or alignment of the spins, magnetization etc.
PC  point contact
PCS  point contact spectroscopy
PSD  power spectral density
RF  radio frequency
RL  read line
RT  room temperature
sec.  section in the manuscript
SD(M)  single domain (model)
SF  spin-flop
SS  scissored state
STT  spin transfer torque
SW  Stoner-Wohlfarth model, switching etc.
TA  transverse
TMR  tunnel magnetoresistance
TMRAM  toggle MRAM
TJ  tunnel junction
V  vortex state or alignment of the spins, magnetization etc.
WL  word line
Chapter 1

Introduction

«The last thing one knows when writing a book
is what to put first.»
Blaise Pascal

One of the emerging branches in today’s microelectronic technology is the spin-based electronics or spintronics. Along with two fundamental physical properties of mass and charge, the electron has a third intrinsic property of spin, which is at the core of spintronic devices. It is difficult to date the exact birthday of Spintronics. Over the last one hundred years the results obtained from many theoretical studies and experimental discoveries in magnetism, solid state physics and quantum mechanics laid the necessary groundwork for spintronics. The most important characteristics of a spintronics device is magnetoresistance (MR). MR is the property of a material to change its electrical resistance under the influence of external magnetic fields. There are several types of MR. The very first experimental observation of MR (anisotropic MR or AMR) was made by W. Thomson in 1856. Tunneling MR (TMR) in magnetic tunnel junctions (MTJ) was observed by M. Julliere in 1975 [1]. Spin-dependent transport in metallic ferromagnetic/non-magnetic/ferromagnetic (F/N/F) structures was first studied in spin injection experiments by M. Johnson and R. H. Silsbee in 1985 [2]. Giant MR (GMR) in alternating F/N layers was discovered by P. Grünberg and A. Fert in 1988-89 [3, 4]. This latter discovery started a huge amount of fundamental and applied research worldwide aimed at using the spin of the electron in microelectronic circuits to improve the existing devices or achieve new functionality. The work, quite deservedly, was awarded the Nobel Prize in 2007.

The experimental discoveries of the spin-dependent transport in metallic N/F multilayers and magnetic tunnel junctions triggered numerous theoretical studies in the end of 1980’s and beginning of 1990’s. The brightest theoretical contributions to the field were made by John Slonczewski, who’s pioneering models of the spin-dependent conductance in MTJ [5, 1989] and current induced spin-
CHAPTER 1. INTRODUCTION

Transfer torque (STT) and spin-wave excitations in magnetic multilayers [6, 7] remain among the most cited works today.

The first commercialized spintronic device was a magnetoresistive readout sensor used in hard drives. This device evolved from AMR to GMR and now to TMR-based sensors, enabling the great increases in the areal data density\(^1\). The technological improvement of the hard drive heads continues. Thus, Hitachi\(^\circledR\) plans to increase today’s highest areal densities to the Tbits/in\(^2\) range by introducing current perpendicular to the plane (CPP)-GMR sensors.

Another important spintronics device is magnetic random access memory (MRAM). The first MRAM research programs started in the 1990’s and today a 4 Mbit capacity, 180 nm feature size MRAM is already commercially available from Freescale Semiconductor\(^\circledR\). Beside the non-volatility and low power consumption of MRAM, it is believed that MRAM will eventually win the race for the operation speed and areal density against the present memory technologies based on conventional electronics\(^2\). Just as in the past twenty years of computer history, the CPU’s smallest feature size evolved from 1 µm in 1985 down to today’s 65 nm. It appears to be only a question of time before the physical dimensions of magnetic memory cells reach ~10 nm size. Such technological achievement requires further theoretical and experimental research in material science, solid state physics and magnetism applied to extremely small structures with dimensions of the order of nanometers.

The topic for this PhD work is to search for experimentally and understand theoretically spin dependent transport at the nanoscale.

The first part of this work focuses on experimental studies of transport through single N/F interfaces in magnetic point contacts (PC). Such PCs can be formed when a sharpened non-magnetic tip is mechanically brought into contact with a ferromagnetic film. The technique allows one to investigate spin dependent transport in very small objects, down to 1 nm in size. Chapter 2 starts with a short theoretical introduction to the method of Point Contact Spectroscopy (PCS) and continues with a description of the PCS measurement technique, concluding with a presentation of the main experimental results. A detailed analysis of the observed magneto-transport is performed. This allowed us to demonstrate new effects with regard to ballistic versus diffusive and surface versus bulk spin transport.

Chapter 3 is a study of the real-time spin dynamics of the coupled magnetic layers in spin-flop (SF) tunnel junctions. SF junctions are used in modern toggle MRAM, making the memory stable against thermal and half-selected switching. The basic principles of the toggle switching as well as the experimental results from the high-speed measurements are presented. The results are compared to analytical single-domain and numerical micromagnetic models. A good agreement between theory and experiment is achieved, demonstrating new spin-dynamic modes in the system.

\(^1\)Usually measured in Gbits/in\(^2\) for hard drive storage media.

\(^2\)Such as FLASH.
Chapter 2

Magnetic Point Contacts

2.1 Theory

In this work point contacts are used as a powerful tool for studying spin dependent phenomena in thin ferromagnetic films when an originally unpolarized current is injected into the films in the CPP geometry. The purpose of this section is to make a short theoretical introduction by presenting the fundamental relationships describing the point contact spectroscopy (PCS) technique as well as for understanding the physics of spin dependent transport in magnetic point contacts. More details can be found in the most complete review on point contact spectroscopy by Naidyuk and Yanson [8].

Despite that there exist several different methods to create a metallic point contact, the simplest way, which has been used in this work, is to mechanically bring in touch to metals. The contact is called a heterocontact when two metals are different. The boundary between these two metals is called an interface. If one of the metals is non-magnetic and the other is ferromagnetic then, such an interface is called non-magnetic/ferromagnetic (N/F) and the contact is magnetic.

2.1.1 Point contact resistance

The basic characteristics of the point contact is its electrical resistance, \( R_{PC} \), and geometrical size, \( d \). As it was noticed by J. C. Maxwell, these two quantities are connected with each other in a simple way:

\[
R_M = \frac{\rho}{d},
\]

(2.1)

where \( \rho \) is the resistivity of the metal\(^1\) and \( d \) is the diameter of the contact orifice, often assumed to be circular. The expression above describes the resistance of a dirty contact in the limit when the size of the PC is much greater than the mean

\(^1\)Generally, \( \rho \) is dependent on temperature.
free path $l$ of the electrons ($d \gg l$), i.e. when electrons scatter many times on the characteristic length scale $d$. The most common scatters are phonons, though it can also be impurities in the metal, as well as magnons in magnetic materials.

### 2.1.2 PC current flow regimes: ballistic, diffusive and thermal

The electrons can experience elastic and inelastic collisions with scattering centers. Depending on the relationship between the PC size, $d$, and corresponding mean free paths, $l_{el}$ and $l_{in}$, the characteristic PC resistance and current flow regime are different. When $l_{el}, l_{in} \gg d$ the regime is ballistic, when $l_{el} \ll d \ll \sqrt{l_{el} l_{in}}$ the regime is diffusive and finally $l_{el}, l_{in} \ll d$ corresponds to a thermal regime. The electron trajectories for each regime are schematically shown in Fig. 2.1. Thus, the transport regime for a given PC depends on the PC size and the amount of impurities in the contact region.

**Ballistic regime**

In 1965 Yu. V. Sharvin theoretically derived the resistance of a ballistic PC [9]:

$$R_{Sh} = \frac{16 \rho l}{3\pi d^2}. \quad (2.2)$$

Expression 2.2 is the so called Sharvin resistance which depends only on the PC size, $d$, and is independent of the material purity. The product

$$\rho l = p_F / n e^2 \quad (2.3)$$

is a constant for a given metal. It is interesting to compare $R_{Sh}$ with the Maxwell resistance given by 2.1. In the limit when $l \simeq d$, $R_{Sh}$ and $R_M$ are comparable.
2.1. THEORY

Diffusive regime

A more general formula for the point contact resistance was obtained by G. Wexler in 1966 [10], [8, pp. 26,31]:

\[ R_{PC} = \frac{16 \rho l}{3 \pi d^2} + \frac{\beta \rho}{d}, \]  
(2.4)

where \( \beta \simeq 1 \) when \( l_{cl} \ll d \). As can be seen from 2.1 and 2.2, the Wexler formula is just a sum of the Maxwell and Sharvin resistances in the ballistic \( (l \gg d) \) and diffusive \( (l \ll d) \) limits. The first term in 2.4 represents the ballistic electron transport and is dominant when the PC size is small, whereas the second term corresponds to the diffusive flow of electrons and is dominant for large contacts and for metals with high resistivity.

Thermal regime

The contact resistance in the thermal regime can be well estimated by the Maxwell formula (2.1). Since the mean free path is much smaller then the contact size, the heat removal from such a PC is not as efficient as in the ballistic or diffusive regimes. The energy dissipation occurs in the contact constriction, leading to a Joule heating as the bias voltage is increased. The temperature of the PC as a function of the applied voltage can be estimated as follows [8, p. 29]:

\[ T_{PC} = \sqrt{T_{bath}^2 + \frac{V^2}{4L}}, \]  
(2.5)

where \( T_{bath} \) is the temperature of the bath (cryostat) and \( L \) is the Lorenz number [11, p. 168].

2.1.3 Point contact spectroscopy

The theory of the point contact spectroscopy has been developed in 1977 by I. O. Kulik. He derived the electron distribution function for a ballistic point contact and showed that an application of voltage \( V \) to a PC results in a splitting of the Fermi surface, with the maximal energy difference equal to \( eV \). Such a splitting on the two sides of a ballistic PC leads to the energy separation between any two electron states at the Fermi surface equal to either zero or \( eV \) (on the same side or on the opposite sides of the PC). This allows to probe the energies of the quasi-particles in the system such as phonons, magnons etc. By varying the voltage across the PC we force the electrons, passing through the constriction to have different initial and final energies. The final electron energy (state) is greatly influenced by the energy-dependent scattering processes, where quasi-particles are involved. This energy-dependence is reflected in the PC resistance as follows.

---

\(^2\)See references at the end of the Ch. 3 in [8].
CHAPTER 2. MAGNETIC POINT CONTACTS

Assuming $\beta \simeq 1$ and $d/l \to 0$ in the Wexler expression 2.4, and using $l = v_F \tau$, we can write:

$$R_{PC} \simeq \frac{16\rho l}{3\pi^2 d^2} + \frac{\rho}{d} = R_{Sh} \left(1 + \frac{3\pi d}{16l}\right) = R_{Sh} \left(1 + \frac{3\pi d}{16v_F \tau}\right),$$

(2.6)

where $\tau$ is the energy-dependent scattering time, which can be related to the electron-phonon interaction function $\alpha^2 F(\epsilon)$ [8, p. 9]:

$$\tau^{-1}(eV) = \frac{2\pi}{\hbar} \int_0^{\epsilon_{\text{F}}} \alpha(e)^2 F(e)\,de.$$  

(2.7)

Here $F(\epsilon)$ is the phonon density of states (DOS) and $\alpha^2$ is the matrix element squared of the electron-phonon interaction. Then the voltage derivative of 2.6, together with 2.7 yields:

$$\frac{dR_{PC}}{dV} \simeq R_{Sh} \frac{3\pi^2 e\alpha^2}{8\hbar v_F} F(eV).$$

(2.8)

The first experimental evidence for 2.8 was obtained by I. K. Yanson in 1974 [12]. It turned out that $\frac{dR_{PC}}{dV} \sim \frac{d^2 V}{dI^2}$, the second derivative of the current-voltage characteristic (IVC), is proportional to the electron-phonon interaction function. As can be seen from Fig. 2.2(a), Yanson identified two distinct maxima in his measured $\frac{d^2 V}{dI^2}$ spectra with the longitudinal (LA) and transverse (TA) acoustic phonon peaks.

The reason why $\frac{d^2 V}{dI^2}$ is proportional to the phonon spectrum is that the electron-phonon interaction results in back-flow scattering in a such way that some of the electrons passing the constriction, are reflected back through the same constriction. These backscattered electrons reduce the net current at energies (bias voltages) corresponding to the phonon modes in the material, and IVC becomes nonlinear. I. O. Kulik theoretically derived the expression for the second derivative:

$$\frac{1}{R_0} \frac{dR(V, T)}{dV} = \frac{8\epsilon d^3}{3\hbar v_F} \int_0^{\infty} \frac{d\omega}{k_B T} g(\omega) \chi \left(\frac{\hbar \omega - eV}{k_B T}\right),$$

(2.9)

where $g(\omega)$ is the point contact interaction function and $\chi(x) = \frac{\alpha^2}{dx^2} \left(\frac{x}{\exp(x) - 1}\right)$ defines the thermal smearing. This thermal smearing often limits the spectroscopic resolution of the PCS technique. $g(\omega)$ is in general a rather complex function. However, at low temperatures and assuming a circular geometry, it can be approximated as $g(\omega) = \alpha^2 F(eV)$ (Eq. 2.8).

PCS is an excellent tool for distinguishing the current flow regimes. This is based on the evaluation of the relative amplitudes of the phonon peaks in the
2.1. THEORY

Figure 2.2: The first experimental observation and identification of the longitudinal and transverse phonon peaks in the \( \frac{d^2V}{dI^2} \) spectra of point contacts to Pb (a), after [12, p. 8]. Schematic illustration of the typical experimental ballistic, diffusive and thermal spectra (b).

measured \( \frac{dV}{dI} \) spectra. Typical shapes of the experimental \( \frac{d^2V}{dI^2} \) curves are schematically shown in Fig. 2.2(b). The presence of pronounced phonon peaks is a clear distinction of the ballistic regime from the two other regimes. This property was important in our statistical analysis of ballistic versus diffusive spin-dependent transport in magnetic PCs (paper I). Since the electron-phonon interaction function is unique for each metal, the experimental spectral pattern (such as shown in Fig. 2.2) can be distinguished for each metal. The characteristic voltages and energies for the TA and LA phonon peaks of the metals used in this work are listed in table 2.1. For the noble metals the TA phonon maximum is frequently present in \( \frac{d^2V}{dI^2} \) spectra, whereas the LA peak corresponding to higher energy phonons is less intensive and is not always observed.

In the case of heterocontacts between two different metals, the contribution from both metals will be present in the PC spectrum. If the effective heterocontact volume is predominantly filled with one metal, then the resulting phonon spectrum will have pronounced peaks of this metal and weak peaks of the other metal. This feature of the PCS method has been explored in paper III, where the mecha-

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<th>Ag</th>
<th>Au</th>
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<td>TA</td>
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<td>10</td>
<td>19</td>
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Table 2.1: Voltages and energies (mV, meV) for TA and LA phonon peaks.
nism of the surface spin-valve effect was investigated for single N/F interfaces.

2.1.4 Spin-dependent transport

2.1.4.1 Spin polarization & magnetoresistance

Spintronics devices would not be possible without materials allowing to establish and manipulate spin-dependent currents in nanostructures. As shown in sections 2.1.1-2.1.3, it is the electrical resistance of a PC, which provides information about the current flow regime, contact size, composition and electron-phonon interaction\(^3\) in the constriction. The resistance of a non-magnetic metal obeys Matthiessen’s rule, stating that at non-zero temperatures there is a contribution to the resistivity coming from electron scattering on thermal phonons [11, pp. 159-162].

In the case of magnetic materials there is an additional source of scattering – magnons. This means that the spin of the conduction electron, being injected into a ferromagnet, can interact with a spin-wave or the localized atomic spins. Such interaction leads to a change of the electron momentum over a characteristic relaxation time, and, therefore, the resistivity is altered. The interaction mechanism relies on a coupling between the conduction \(s\)-electrons and the \(d\)-electrons in the ferromagnetic material. The latter electrons are responsible for the magnetism itself. As can be seen in Fig. 2.3(a), the \(s\)-band typically possesses a symmetrical density of states (DOS) distribution, \(D^{\uparrow \downarrow}(E)\) with respect to spin-up (\(\uparrow\)) and spin-down (\(\downarrow\)) electrons. In contrast, the \(d\)-electrons have asymmetrical DOS, where one

\[D^{\uparrow}(E)\]

\[D^{\downarrow}(E)\]

\(E_F\)

(a) schematic DOS for \(s\)- and \(d\)-bands

(b) calculated bands of 3\(d\)-transition metals: Fe, Co, Ni, Cu (after [13] p. 28)

Figure 2.3: Spin polarization and DOS in ferromagnets.

\(^3\)Generally speaking, electron-quasiparticle interaction.
sub-band is shifted with respect to the other by the amount of the exchange energy, \( E_{\text{ex}} \). In ferromagnetic metal bands corresponding to the different spin orientations are labeled according to their relative spin population. Thus \( d \) sub-band containing the largest number of electrons is called \textit{majority band} and the other band is \textit{minority band}. The magnetization \( \mathbf{M} \) in the material is determined by the minority spins. The majority and minority bands are shown in Fig. 2.3(a) as \( d^\uparrow \) (left) and \( d^\downarrow \) (right) respectively. That is, for each spin channel the number of electrons participating in the conduction is different and we can refer to this difference as \textit{spin polarization}. An example of the calculated spin-polarized bands for Fe, Ni, Co and the unpolarized Cu band is shown in Fig. 2.3(b).

The actual spin polarization process of an unpolarized current passing through a N/F interface is the following. An equal amount of conduction electrons from the \( s^\downarrow \) and \( s^\uparrow \) bands of the N metal enters the ferromagnet. The \( s^\uparrow \) electrons can be scattered into the \( d^\uparrow \) band, while \( s^\downarrow \) electrons cannot be scattered into the \( d^\downarrow \) band due to a lack of free states there. Thus, in ferromagnets the electrical resistivity, \( \rho \), can be ascribed to the resistivities of the \( \downarrow \)- and \( \uparrow \)-spin conducting channels connected in parallel:

\[
\rho = \frac{\rho^\downarrow \rho^\uparrow}{\rho^\downarrow + \rho^\uparrow},
\]

where conductivity \( \sigma^\downarrow, \sigma^\uparrow = \frac{1}{\rho^\downarrow}, \frac{1}{\rho^\uparrow} = \frac{e^2}{\tau^\downarrow, \tau^\uparrow} \frac{D^\downarrow, \uparrow}{m^*} \) for an effective mass \( m^* \), spin-dependent relaxation time \( \tau^\downarrow, \tau^\uparrow \), and DOS \( D^\downarrow, \uparrow \). This model is known as the \textit{two-channel model} and it neglects spin-flip processes (no spin-channel mixing).

In a more realistic situation, when electrons scatter additionally on magnons at high temperatures, expression 2.10 is somewhat modified [14]:

\[
\rho = \frac{\rho^\downarrow \rho^\downarrow \uparrow + \rho^\downarrow \uparrow \rho^\downarrow}{\rho^\downarrow + \rho^\downarrow \uparrow + 4\rho^\downarrow \uparrow},
\]

where \( \rho^\downarrow \uparrow \) represents spin-flip scattering. Naturally, in the limit where \( \rho^\downarrow \uparrow \ll \rho^\downarrow, \rho^\downarrow \uparrow \), expression 2.11 becomes the same as 2.10. The two-channel model of the resistivity for the N/F multilayers was developed by Valet and Fert where authors take into account volume and interface spin-dependent scattering [15].

In the two-channel model, assuming the total electrical current is \( I_{\text{tot}} = I^\uparrow + I^\downarrow \) and the difference between the spin-\( \downarrow, \uparrow \) currents or the spin current is \( I_{\text{diff}} = I^\uparrow - I^\downarrow \), the definition of polarization \( P \) is:

\[
P = \frac{I_{\text{diff}}}{I_{\text{tot}}} = \frac{\sigma^\downarrow - \sigma^\uparrow}{\sigma^\downarrow + \sigma^\uparrow} = \frac{D^\downarrow (E_F) - D^\uparrow (E_F)}{D^\downarrow (E_F) + D^\uparrow (E_F)}. \tag{2.12}
\]

Eq. 2.12 does not take into account the fact that the electronic DOS is typically spin-asymmetric for heavy \( d \)-electrons, while the electrical current is mostly carried by light \( s \)-electrons. A more comprehensive definition of \( P \) is given in Ref. [16]. From the practical point of view it is enough to know that for typical bulk ferromagnetic materials used in the experiments, such as Ni, Co and Permalloy (Ni_{80}Fe_{20}), the
degree of polarization lies within the range of 25–45% [17–19]. It is also known that the polarization in thin films may differ from that in the bulk [20, 21].

The physics of spin polarization discussed above is the foundation for spintronics. An unpolarized current becomes polarized inside a ferromagnetic material. When such current leaves the ferromagnet, e.g., in a N/F/N trilayer, the spin polarization remains even in the non-magnetic material. This phenomenon is known as spin accumulation. Spin accumulation is characterized by its size, which is determined by the equilibrium between the net spin-injection rate at the F/N interface and the spin-flip rate in N. The spin accumulation decays exponentially away from the F/N interface within a so called spin-flip or spin-diffusion length. The spin-flip length depends on the spin-flip rate, $\tau_{\uparrow\downarrow}$, Fermi velocity, $v_F$, and mean free path of the electrons, $l$, as $l_{sf} = \sqrt{\tau_{\uparrow\downarrow} v_F l} / 3$. Now, the question is, what will happen if we place another ferromagnet following the first F/N interface, within the spin-flip length from the interface? Assuming either parallel (P) or antiparallel (AP) orientation of two ferromagnets, as schematically shown in Fig. 2.4(a), the conduction electrons, whose spin orientation is different from the magnetization of the traversed ferromagnetic layer, are scattered more strongly. In other words,

![Figure 2.4: Two-current model of GMR effect in F/N/F multilayers for P and AP magnetization configurations. Electron trajectories (a) and equivalent electrical circuit (b).](image)

spin-$\uparrow$ and spin-$\downarrow$ electrons experience different total resistances (equivalent circuit is shown in Fig. 2.4(b)). Such a F/N/F trilayer or a spin-valve is characterized by the magnetoresistance (MR) ratio:

$$MR = \frac{R_{AP} - R_P}{R_P},$$

(2.13)

where $R_P$ and $R_{AP}$ are the resistances in the parallel and antiparallel alignments of the magnetizations of the ferromagnetic layers.
2.1. THEORY

GMR  In the popular literature as well as in scientific papers it is common to refer to MR as giant magnetoresistance (GMR) in the case when MR originate from the N/F alternating layers.

In the general case, when the magnetizations of the F layers are aligned at an arbitrary angle \( \theta \), the total resistance of the spin valve can be expressed as follows:

\[
R = R_p + \Delta R (1 - \cos \theta) / 2 = R_p + \Delta R \sin^2 \frac{\theta}{2},
\]

where \( \Delta R = R_{AP} - R_P \).

AMR  The first observed MR was the anisotropic magnetoresistance, W. Thomson in 1856. The experimental observation was that the resistivity of a ferromagnetic material is a function of the angle \( \theta \) between the magnetization direction and the electrical current. In the normal case, the resistivity is maximum \( (\rho_\parallel) \) when the current is parallel to the magnetization direction, and minimum \( (\rho_\perp) \) for perpendicular orientation:

\[
\rho(\theta) = \rho_\perp + (\rho_\parallel - \rho_\perp) \cos^2 \theta = \rho_\perp \sin^2 \theta + \rho_\parallel \cos^2 \theta.
\]

The origin of the AMR effect is in the anisotropy of scattering produced by spin-orbit interaction [22, 23]. The typical AMR value is 0.1–3% for pure metals and can be somewhat higher for alloys. It is known that for Co films, used in much of this work, the AMR is about 0.8% according to [24].

TMR  If the metallic N spacer, shown in Fig. 2.4(a), is replaced by an insulator (I) the structure becomes a F/I/F magnetic tunnel junction (MTJ). Though MTJs were not used in the PC measurements presented in this chapter we mention the main relationship for the MTJ MR (Eq. 2.17), which is used in the spin-flop dynamics experiment discussed in Ch. 3.

In 1975 M. Julliere discovered tunnel magnetoresistance (TMR) in Fe/Ge/Co tri-layers [1]. Though the TMR can be estimated using 2.13, Julliere derived a simple expression for the TMR, where he assumed no spin-flip tunneling of the conduction electrons and that the number of electrons in each F layer is proportional to the DOS at the Fermi level:

\[
\text{TMR} = \frac{2P_1 P_2}{1 - P_1 P_2},
\]

where \( P_{1,2} \) are the polarizations of the ferromagnetic electrodes. The disadvantage of the Julliere model is that neither height nor thickness of the tunnel barrier is taken into account. A more accurate theoretical model of the MTJ was presented by J. Slonczewski in 1989. He solved the Schrödinger equation for the single electron Hamiltonian in the case of the rectangular tunnel barrier separating two free-electron like ferromagnetic metals. As the main result he derived the resistance

\[\text{\footnote{Originally the conductance has been derived, but for practical use we give an equivalent expression for the resistance of the MTJ.}}\]
of the MTJ as function of the angle $\theta$ between the magnetization vectors of the two ferromagnets:

$$R(\theta) = \frac{R_0}{1 + P^2 \cos \theta},$$

(2.17)

where $P$ is the effective spin polarization of the ferromagnetic-barrier couple:

$$P = \frac{(k_{\uparrow} - k_{\downarrow})}{(k_{\uparrow} + k_{\downarrow})} \left( \frac{\kappa^2 - k_{\uparrow} k_{\downarrow}}{\kappa^2 + k_{\uparrow} k_{\downarrow}} \right).$$

(2.18)

Here $k_{\uparrow}, k_{\downarrow}$ are the Fermi wave vectors and $\kappa = \sqrt{(2m/\hbar^2)(V - E_F)}$ is the constant which depends on the tunnel barrier height $V$.

Slonczewski’s model is a good approximation of the magnetconductance in real MTJs. A comparison of his and Julliere’s models with numerical calculations performed by J. M. MacLaren et al. in [25].

**Domain wall magnetoresistance (DWMR)** A domain wall (DW) is an object connecting regions of a magnet in which the magnetization has a different orientation, as shown in Fig. 2.5 where the magnetization on the right and left sides of the sample is pointing along the same axis but in opposite directions. The smooth connection between the two regions is due to a competition between the exchange energy, $E_{\text{ex}}$, and the anisotropy energy, $K$. $E_{\text{ex}}$ works to keep $\Delta \theta$ small, whereas $K$ works to have as few spins as possible in the direction perpendicular to the anisotropy axis. Thus, the DW is formed with a non-zero width which can be estimated [26, p. 331] as $w = a\sqrt{E_{\text{ex}}/2K}$ or, equivalently, $w = 2\sqrt{A_{\text{ex}}/K}$. Here $a$ is the lattice constant of the material and $A_{\text{ex}}$ is the exchange stiffness.

DWs have a high application potential for information storage. It has been shown experimentally [27–30] that the critical current densities required to move DWs can be significantly smaller than those required for the current induced magnetization reversal (CIMR). The domain wall magnetoresistance (DWMR or DMR),
similar to the GMR effect, arises from the fact that the polarized conduction electrons traversing the DW experience spin-dependent scattering by the magnetization rotating along the electrons trajectory, which leads to an increase in the resistivity. As an example, a ferromagnetic nano-wire having at least one DW will have a higher resistance than a saturated single domain wire. The estimation of the DWMR is a rather non-trivial task and is out of scope for this experimental work. However, studies have been performed where DWMR was estimated for nano-wires, thin films, magnetic point contacts, etc. [31–36]. It is worth to also mention theoretical [32, 37–39] as well as recent experimental [40, 41] studies on current-driven dynamics of DW’s, which are interesting for technological applications demanding high-speed read/write access to the magnetic storage medium. Current-driven DW dynamics relies on the mechanism of spin momentum transfer in a current carrying ferromagnet. The basics of the spin momentum transfer are discussed in the next section.

One way to create a stable (stationary) DW in a nano-wire is to make a geometrical constriction, such as a notch or a bending in the wire. Then a DW can be trapped by the pinning potential of the constriction. Our micromagnetic simulations show, however (section 2.3.2, paper IV), that in thin ferromagnetic films the DW can be pinned by the anisotropy energy alone, without any geometrical constrictions.

2.1.4.2 Spin transfer torques

Spin transfer torques (STT) have been intensively explored during the last 20 years. In the middle of the 80s, theoretical and experimental studies on current-induced DW motion in ferromagnetic films [42–44], showed that an electrical current sent through ferromagnetic films or wires containing DW, interacts with the walls and, if the current density is high enough, can cause DW motion. In 1989 Slonczewski calculated STT for the case of current traversing on F/I/F MTJ [5]. It turned out to be impractical to apply torques across TJs, and the following theoretical works [6, 45] by Berger and Slonczewski predicted much more efficient spin transfer in metallic multilayers, such as F/N/F. These papers triggered further theoretical as well experimental research [46–52] on STT.

In his early work Slonczewski calculated analytically the STT as a transverse exchange component [5, Eq. 5.4] for a F/I/F multilayer, and later incorporated STT [6, 51] into the Landau-Lifshitz-Gilbert (LLG) equation for the magnetization dynamics as an additional term:

\[
\frac{d\mathbf{m}}{dt} = -|\gamma| \mathbf{m} \times \mathbf{H}_{\text{eff}} + \alpha \mathbf{m} \times \frac{d\mathbf{m}}{dt} + \tau(I, \mathbf{m}, \mathbf{P}),
\]

(2.19)
where $\alpha$ is damping parameter, $H_{\text{eff}}$ is an effective field which is a sum of the external, demagnetizing, intrinsic anisotropy and fields present in $F$. $\tau(J, m, P)$ is the STT density and $m = M/M_s$ is the reduced magnetization of the F layer.

The LLG equation (2.19) describes the temporal evolution of a magnetic moment $m$ in a ferromagnet when a non-zero magnetic field is applied and in the presence of dissipation characterized by the damping parameter $\alpha$. In the macrospin approximation it is assumed that magnitude $|m|$ remains constant in time. A detailed review of the precessional magnetization dynamics described by the LLG equation can be found in the book on magnetism by Stöhr and Siegmann [53, Ch. 3].

Consider a $F_p/N/F$ device$^5$ (spin-valve) as shown in Fig. 2.6, where the electrons flow from the left to the right (negative current) traversing the two ferromagnetic layers with the magnetizations oriented in the plane of the layers. We assume that $F$, referred as free layer, is thinner than $F_p$, referred as fixed layer, and is easier$^6$ to reorient by STT. The fixed layer serves as a polarizer for the unpolarized electrons and, assuming that the thickness of the spacer $N$ is smaller than the spin diffusion length in $N$, the conduction electrons are spin-polarized when entering the free layer. If the magnetic moment of the free layer is oriented at an angle to the moment of the fixed layer, the free layer will absorb the component of the incoming spin-polarized current transverse to the layer’s magnetic moment, as a result

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{figure2_6.png}
\caption{The STT effect in an $F_p/N/F$ spin-valve structure. Electronic current is polarized by $F_p$ and exerts STT on $F$. The spin precession in $F$ about $H_{\text{eff}}$ is described in the text and is illustrated in more detail in Fig. 2.7(right).}
\end{figure}

\footnote{Strictly speaking $N/F_p/N/F/N$ multilayered device.}

\footnote{The fixed layer can be made of much stiffer magnetic material or can be \textit{pinned} to another layer, magnetically very hard layer.}
the free layer will feel a torque tending to rotate its moment toward the orientation of the incoming spins, as schematically illustrated in Fig. 2.7(left). The spin torque vector will act on the magnetic moment of the free layer in the opposite direction from the magnetic damping torque vector, as shown in Fig. 2.7(left) by the black dashed arrows. Mathematically, all participating torques are depicted in equation 2.19 (and Eq. 2.20). Even intuitively, we can see from 2.19 that the motion of the magnetic moment can be circular or spiral about the axis parallel to \( H_{\text{eff}} \) due to the fact that all three torques are orthogonal to each other.

The remarkable property of the F\(_p\)/N/F multilayer confined into a nanopillar geometry, such as shown in Fig. 2.6, is that the magnetization of the free layer can be switched by the current along the easy axis. The electrons flowing from the fixed layer become polarized and exert a torque on the free layer, and at a high enough current this leads to a magnetization reversal of the free layer. The result is that the magnetizations in the two layers are parallel (P state) and total resistance of the pillar is low (positive current range, as illustrated in Fig. 2.8). When the conduction electrons flow in the opposite direction, from the free to the fixed layer, the switching mechanism is different. Spin-dependent scattering at the F\(_p\)/N interface leads to a spin accumulation in N, where reflected electrons are polarized parallel (AP) to F\(_p\). The electrons with spins parallel to the magnetization in F\(_p\) are preferentially transmitted through F\(_p\). Thus, backscattered current is spin-polarized, but is AP to F\(_p\). If the current is high enough, then this spin accumulation will create a net torque on the moment of the free layer which can lead a switching of the latter into the AP state of the valve. The typical dependence of resistance on the bias current for a pillar device is shown in Fig. 2.8. A detailed study of such hysteretic CIMS as well as thermally activated switching in F\(_p\)/N/F trilayers was performed...
Figure 2.8: Illustration of the typical hysteretic \( R(I) \) dependence for the \( F_p/N/F \) trilayer nanopillar device shown in Fig. 2.6. Horizontal arrows schematically indicate direction of the electron flow, whereas vertical arrows \( \uparrow \) and \( \downarrow \) (\( \uparrow \downarrow \)) shows relative magnetization orientation of the polarizing \( F_p \) and sensing \( F \) ferromagnetic layers respectively. The \( R \) dependence on externally applied magnetic field is shown in the inset.

by Albert et al. [55] and can be consulted for more details.

The theory of Slonczewski of spin-transfer torques has been validated by J. Xiao et al. [56] using more detailed numerical calculations. The expression for the STT density term \( \tau(J, m, P) \) in the LLG equation 2.19 can be written in the form:

\[
\tau(J, m, P) = |\gamma| \beta \epsilon (m \times m_p \times m),
\]

where the following notations are used:

- \( \gamma \) – Gilbert gyromagnetic ratio,
- \( m_p \) – unit vector of electron polarization direction,
- \( \beta = \frac{\hbar}{\mu_0 e |tM_s'|} \)
- \( e \) – electron charge,
- \( t \) – thickness of magnetic layer,
- \( \epsilon = \frac{PA^2}{(\Lambda^2 + 1) + (\Lambda^2 - 1)(m \cdot m_p)} \)
- \( \Lambda^2 = \frac{1}{2}(R_\uparrow + R_\downarrow)G \)
- \( G \) – conductance of \( N \).

Magnetization precession and CIMR become pronounced when the STT term 2.20 becomes comparable with the damping torque in 2.19. Naturally, this depends on
the current density $J$. Strong STT (high $J$) is achieved in spin-valves with small cross sections. As reviewed in [54, 57] and illustrated in Fig. 2.7(right) for the macrospin approximation, the magnetic moment can perform three types of precessional motion when subject to spin torque: when the current is below a certain threshold value the excited magnetic moment relaxes back to the easy axis; if the current is just above the threshold then the moment makes many precessional cycles before it reverses its direction; and finally, much higher current than the threshold value leads to a quick magnetization reversal.

The first direct electrical measurements of the microwave-frequency dynamics driven by the polarized current in nanomagnets has were performed by S. Kiselev et al. [58]. We will consider the topic of field induced magnetization dynamics in chapter 3. In the following sections we concentrate the reader’s attention on the magnetic excitations caused by spin transfer torques.

2.1.4.3 Current-driven excitations in N/F multilayers

The first experimental observation of current-driven magnetic excitations as well as the first indirect evidences of RF magnetization dynamics have been reported by M. Tsoi et al. [47, 59] in point contacts to Cu/Co multilayers. The measured IVC ($dV/dI$) as a function of bias voltage and externally applied magnetic field $H$ (Fig. 2.9(a)) revealed distinctive peaks in $dV/dI(V)$, which, due to their in $V$ dependence on $H$, were attributed to an excitation of spin waves. This work was an important confirmation of the Berger and Slonczewski theoretical predictions [6, 45] and the following observations were made: (i) the measured excitations were on the negative side of the IVC, corresponding to the case when the electrons flow from the non-magnetic tip into the multilayer; (ii) the excitations are upward.
steps or peaks in the resistance; (iii) the position of the peaks in bias varies linearly with $H$; (iv) the importance of the diffusive nature of the electron transport in the PCs was noted. The latter observation will be discussed in more detail in section 2.3.1 (paper I).

### 2.1.4.4 Current-driven excitations in single F layers

The main subject of this PhD work is spin-dependent transport phenomena in PCs to single ferromagnetic films. Thus, in this section we review the experimental and theoretical studies on STT in single ferromagnetic layers performed prior to our work.

Already in 1999 the Cornell University group (E. B. Myers et al.) [49] studied CIMS in N/F multilayers found magnetic excitations in IVCs for a single N/F interface as shown in Fig. 2.10. A more detailed experimental study on PCs to a single F layer [60] revealed strong similarities between the current-induced magnetic excitations in N/F multilayers and in single ferromagnetic layers, as shown in Fig. 2.9. Since there is no spin polarizing ferromagnetic layer in the single F-
2.1. THEORY

layer case, the main question is what mechanism is responsible for the magnetic excitations caused by the unpolarized current?

M. L. Polianski and P. W. Brouwer in their theoretical work [61] proposed that an unpolarized current traversing a single thin F layer can create a spin-wave instability, which occurs for one current direction only and depends on the asymmetry of the N/F/N structure. The mechanism of this instability is the same as the one causing the AP CIMR in F/N/F multilayers described above, i.e., the STT arising from the spin accumulation in the normal metal spacer near the N/F interface. The asymmetry of the N₁/F/N₂ structure means that the spin accumulation is different at the N₁/F and F/N₂ interfaces due to differences in geometry, impurity concentration, etc. This model has been successfully verified by B. Özyilmaz et al. [62], where symmetric pillar devices showed no magnetic features in their IVC and asymmetric pillars showed magnetic excitations. Independently from the Polianski and Brouwer work, M. D. Stiles et al. [63] theoretically showed that in the case when an unpolarized current traverses a symmetric N/F/N device, precession-type spin-wave instabilities in F are possible if the magnetization in F is non-uniform, as illustrated in Fig. 2.11. The electron having an AP spin align-

![Figure 2.11: Illustration of the STT at the N/F interface due to a magnetization non-uniformity in F and impurity scattering in N of originally unpolarized electrons, after [63].](image)

ment with the local magnetization M₁ in F (not shown in the figure) experiences back scattering into N material. Due to impurities present in N near the N/F interface this electron diffuses into the ferromagnet again where local magnetization M₂ can be different from M₁. Such difference leads to an absorption of the trans-
verse spin component of the electron and an appearance of the net STT acting on local spins (Nst, Fig. 2.11). The magnetization non-uniformity along the current path is itself a symmetry breaking condition, which can be sufficient for the appearance of the magnetic excitations for single N/F interfaces. Our experimental verification of the Polianski-Brouwer and Stiles et al. theories on single-interface spin transfer torques is discussed in detail in Sec. 2.3.1 (paper I).

Rather surprisingly, researchers observed even hysteretic switching in PCs to single ferromagnetic films [64], similarly to spin-valve trilayers. This effect is the subject of the discussion in section 2.3.2 (papers II, III, IV and [65]).

2.2 Experimental techniques

2.2.1 Sample fabrication

In this section a short description of the sample fabrication is given. Metallic multilayered films were used in experiments with magnetic point contacts. Such multilayers were created using e-beam evaporation. In our case, an ultra high vacuum deposition system from Eurovac® with an e-gun from Thermionics® was used.

The films were deposited at room temperature and typical $10^{-7}$ mBar base pressure in the vacuum chamber. 1 mm thick silicon wafers coated with 1 µm of SiO$_2$ were used as substrates (Fig. 2.12(a)). The first deposited metal is a thin layer of Ti serving as a buffer layer, followed by a relatively thick layer of noble metal, such as copper, used as the bottom electrode in point contact measurements. Finally, a ferromagnetic film was deposited with an optional protective coating, as shown in Fig. 2.12(b). This protective layer of Au or Cu prevents the ferromagnetic layer from oxidizing when sample is subjected to air. The typical fabrication recipe is

![Figure 2.12: Typical fabrication process (not to scale)](image-url)
shown in table 2.2. We deposit thin metallic layers at \(\lesssim 1 \text{ Å/s}\) in order to minimize the roughness of the films.

<table>
<thead>
<tr>
<th>step</th>
<th>deposited metal</th>
<th>thickness (nm)</th>
<th>rate (Å/s)</th>
</tr>
</thead>
<tbody>
<tr>
<td>1</td>
<td>Ti</td>
<td>2…10</td>
<td>(\sim 1)</td>
</tr>
<tr>
<td>2</td>
<td>Cu</td>
<td>20…100</td>
<td>(\lesssim 5)</td>
</tr>
<tr>
<td>3</td>
<td>Co</td>
<td>2…100</td>
<td>(\lesssim 1)</td>
</tr>
<tr>
<td>(optional) 4</td>
<td>Au</td>
<td>2…5</td>
<td>(\leq 1)</td>
</tr>
</tbody>
</table>
2.2.2 Measurement technique

«One thing I have learned in a long life:
that all our science, measured against reality,
is primitive and childlike – and yet
it is most precious thing we have.»

Albert Einstein

2.2.2.1 Cryogenic setup

All experiments with magnetic point contacts have been conducted at liquid Helium (4.2 K) temperatures. The experimental setup, schematically shown in Fig. 2.13, consists of a vacuum insulated reservoir (or dewar) where liquid He is stored, a cryostat, a superconducting solenoid magnet, and a sample holding mechanism (or fixture).

The key feature in this setup is the fixture with a built-in mechanism allowing to establish point contacts between a metallic tip and the sample. The differential

Figure 2.13: Schematic of the cryogenic experimental setup used in PC measurements.
screws inside the fixture cylinder provide the necessary XZ positioning control of
the metallic tip over the sample. Since the typical sample dimensions are about
10 × 5 mm, the $X_{\text{ctrl}}$ allows to change the position of the tip along the sample,
whereas $Z_{\text{ctrl}}$ moves the tip up or down with respect to the sample plain.

All electrical wires used for 4-point measurements (see Sec. 2.2.2.2), the tem-
perature sensor and the coaxial cable are inside the cylinder of the fixture and can
be accessed through the connector at the fixture top.

2.2.2.2 Electrical measurements

The magnetoresistive characterization of the metallic point contacts to ferro-
magnets as well as point contact spectroscopy requires sensitive electrical mea-
surements of the current-voltage characteristics (sec. 2.1.3). To achieve a high de-
gree of sensitivity the so-called modulation technique was used. Though this tech-
nique is reviewed at great extend in [8, Ch. 4.3] here we give a short description of
the electrical measurement setup used in this work.

The electrical circuit used in point contact measurements is shown in Fig. 2.14.
Voltage $V_0$ is supplied by a function generator such as Agilent® 33120A, which is
can output currents up to several tens of mA. Voltage $V_s$ and current $I_s$ through
the sample are measured by DC multimeters in a standard 4-point measurement
configuration. Since all measurements were performed in the current bias regime,
where bias resistor $R_b$ is used to limit and stabilize the current flowing through
the point contact. In addition to the $V(I)$ of the point contact we measure its
first, $V_1 \sim dV(I)/dI$, and second, $V_2 \sim d^2V(I)/dI^2$, derivatives using Standfor
Research® SR830 lock-in amplifiers. To measure the IVC for a given current range
the function generator is programmed to output triangular voltage waveform at
low 0.5 – 3 mHz frequency.

The AC output of the first lock-in amplifier is controlled by the $R_b^{AC}$ resistor
and typically is much smaller than the DC bias voltage applied to the point con-
tact. The frequency of this AC signal is fixed to $\omega/2\pi \approx 483$ Hz. This signal is
mixed with the DC bias using an LC circuit.

$V_2$ ($2^{nd}$ lock-in amplifier) is required for phonon spectroscopy measurements
(see sec. 2.1.3). An LC bandpass filter is used to measure $V_2$ as the second har-
monic ($2\omega$) of the point contact response to the AC signal fed by the $1^{st}$ lock-in
amplifier.

The magnetic field $B$ is produced by the solenoidal magnet powered by a
Kepco® bipolar power supply. The magnet is capable of producing upto 8 T mag-
netic field. The field is applied in the plane of the sample. The exact value of the
field is measured as a voltage over 0.1 Ω resistor standard connected in series to
the magnet. The superconducting solenoid calibration constant is 0.22 T/A.

All equipment mentioned above was connected to a computer via GPIB. The
measured data were read out and stored by the computer software designed in the
LabView® programming environment.
2.3 Results and discussions

2.3.1 Phonon spectroscopy and spin transfer torque in magnetic point contacts (paper I)

As discussed in Sec. 2.1.4.4, theoretical predictions of current-induced magnetization excitations were made also for single ferromagnetic layers [61, 63]. The
mechanism of such excitations relies on spin transfer in the direction normal to the current flow. As illustrated in Fig. 2.15, a spin-dependent reflection of an electron can be viewed as a transfer of a magnetic moment from the F layer to the back-scattered electron, which thereby becomes spin-polarized near N/F interface (electron 2 with dashed trajectory). As will be shown below, an impurity in N plays the role of a scatterer, which can reflect the electron back to the F layer. This scattering leads to a spin-diffusion process, in which a spin transfer of the magnetic moment occurs between points A and B of the interface. The spin transfer is mediated by the transport electrons, if the current is strong enough and the back-scattering rate is high, a net spin torque is produced that can excite spin waves in F the direction normal to the current flow [61]. Such single N/F interface STT has been used to interpret excitations in the differential resistance of magnetic PCs and nanopillars having single F layer [60, 62].

A direct experimental verification of this new mechanism of spin torque, induced by back-scattered electrons, would probe the strength of the impurity scattering near the N/F interface and correlate it with the magnetic excitations. Such study has been performed in this paper by using PCS. As briefly reviewed in section 2.1.3 (p. 7), the main advantages of using PCS are the possibility to distinguish the regime of the current flow in the PC and at the same time detect magnetic excitations by measuring the resistance of the PC. The ballistic regime can be distin-

Figure 2.15: Schematic representation of the single N/F interface STT mechanism. The electron (0) with spin parallel to the film’s magnetization experiences no reflection at N/F interface. Depending on orientation of the electrons (1,2) spin with respect to magnetization in M, impurities in N, and presence of the magnetization nonuniformity in F near N/F interface, the transfer of the spin momentum can occur in F, for example, between points A and B for electron 2 (dashed trajectory).
guished from the diffusive and thermal regimes by analyzing the second derivative \( \frac{d^2V}{dI^2} \) of the measured IVC, or the PC spectrum. As schematically shown in figure 2.2(b), \( \frac{d^2V}{dI^2} \) has pronounced phonon peaks in the ballistic regime, which vanish in the other two regimes. It is interesting to compare PCs having different current flow regimes but nearly the same resistance\(^7\) and therefore the same current density in the nanoconstriction, as shown in Fig. 2.16. Contacts to

![Figure 2.16: A comparison of the PC spectra for pairs of contacts of similar resistance but different current flow regimes: bulk Co and Ag tip measured at zero field (a, b); 100 nm thick Co film and Ag tip (solid line, \( B = 3 \) T) and Cu tip (dotted line, \( B = 4 \) T) (c). The ballistic and diffusive contacts are represented by dotted and solid lines respectively. The threshold current density for the magnetic peaks is \( \sim 10^9 \) A/cm\(^2\).

both bulk and film Co show N-shaped magnetic peaks, which are present only in the diffusive PCs, in the bias range of \(-90 \ldots -50\) mV (solid curves). In contrast, ballistic contacts showed no magnetic excitations (dotted curves). This observation suggests that the current-induced magnetic excitations are determined by the strength of the impurity scattering near the N/F interface. The experiments were repeated for bulk Co, Co thick films as well as thin films (as shown in paper II).

The importance of the impurity scattering in the single N/F interface STT effect is supported by the statistical analysis of the measured point contacts. A spectral

\(^7\)As well as the same PC size estimated by Eq. 2.2 and 2.4.
quality, or degree of the diffusivity of each PC, can be characterized by a phenomenological parameter $\gamma = B/A$, which is inversely proportional to the strength of the phonon peaks $A$ relative to the background $B$, as shown in inset to Fig. 2.17. The PC statistics in 2.17 shows that the probability of observing current-induced

![Figure 2.17: Statistics of the measured PCs with and without magnetic excitations.](image)

magnetization excitations for ballistic PCs ($\gamma < 1$, open squares in figure) is significantly lower than for diffusive or thermal ($\gamma > 1$, filled squares) point contacts. Thus, diffusive regime favors the STT effect for a single N/F interface. Among many measured PCs, not a single highly ballistic contact that would show magnetic anomalies was observed.

It should be noted however that the diffusive character of the current flow may not be a sufficient condition for the occurrence of magnetic excitations. As seen from the figure, some diffusive contacts exhibit no magnetic peaks, which may be the result of some unfavorable spin distribution in the ferromagnet. As proposed in [63], a non-uniform spin distribution is required for observing magnetic excitations at single N/F interfaces (see Fig. 2.11).

The experimental results described above provide a direct confirmation of the proposed single N/F interface mechanism of the current-induced STT [61, 63], which relies on the impurity mediated spin polarization of the originally unpolarized current.

### 2.3.2 Hysteresis in point contacts to single ferromagnetic films

From Sec. 2.1.4.4 we have learned that, similar to magnetic multilayers, single N/F interfaces can exhibit current-induced magnetization excitations, manifested as peaks in resistance when an unpolarized current flows into F. The fundamental origin of such magnetization instabilities is the difference in conductivity between the majority and minority spin channels in the ferromagnet, which leads
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to spin accumulation and therefore magnetic STT at the interface [61, 63]. As has been shown experimentally in Sec. 2.3.1 (paper I), this mechanism relies on strong impurity scattering near the N/F interface, which is characteristic of diffusive or thermal PCs.

The new type of the STT effect described above cannot explain the origin of current-induced hysteretic switching (CIS), which was recently observed [64] in single F films (Fig. 2.18(a)). Qualitatively, this switching is very similar to that in spin-valve devices, F/N/F, where it is well known that \( R(I) \) dependence exhibits hysteresis, illustrated in Fig. 2.8 (p. 18). A natural question arises when looking at this peculiar similarity: *is it possible to have a current-controlled magnetic memory device based on a single ferromagnet only?* This question motivated our research where we tried to clarify the origin of the CIS in single F film. We discuss two possible origins of this hysteresis:

- stress induced anisotropy in the PC core;
- energetically distinct surface spin states, including spin vortex states.

A careful analysis of the experimental data is carried out to verify the applicability of each model.

* * *

Before presenting our modelling of the hysteretic switching for single F layers, we review the first proposal to explain the effect by T. Y. Chen, Y. Ji, C. L. Chien and M. D. Stiles [64] in 2004. The authors proposed the STT effect in a single exchange-biased Co layer 400 nm in thickness. The exchange bias arises from a thin antiferromagnetic (AF) CoO layer formed on the top surface of the Co film due to natural oxidation. In order to create a strong exchange between the AF layer and the top Co surface, the samples in [64] were cooled in a strong magnetic field applied in the film plane. The sign of the initial cooling field establishes the preferential alignment of the top Co surface towards the \(+H\) direction. Thereby a *unidirectional* exchange anisotropy is created. To study STT in such exchange-biased Co film, the authors use the point contact technique, which provides high current densities required for strong spin torques. Measurements of the resistance \((V/I)\) as well as differential resistance \((dV/dI)\) revealed hysteretic CIS, shown in Fig. 2.18(a). A distinctive bistable characteristic of these hysteretic loops suggests that there are at least two magnetic regions, where one switches its magnetization with respect to the other. Thus, the authors consider three magnetic regions as illustrated in Fig. 2.18(a)(inset): *(i)* the bulk of the film, which is unaffected by the exchange bias and current; *(ii)* the near surface region (not shown in inset) which is influenced by the exchange bias but not the current and *(iii)* the *nanodomain* in the contact core. This nanodomain is assumed to be coupled to the AF surface layer and have *two* stable magnetization directions relative to the Co film. Current injected from the tip with high density can exert STT sufficient to switch this nanodomain. We recall, however, that exchange anisotropy is unidirectional and...
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[Figure 2.18: Hysteretic current-induced switching loop at zero field in a Co/Cu(tip) PC (a) and the proposed micromagnetic model (inset), after T. Y. Chen et al. [64]. Illustration of the stable and unstable micromagnetic configurations of the surface antiferromagnetic (AF) layer and the nanodomain (ND) (b).]

therefore there must be only one preferable magnetic configuration of the ND, as shown in Fig. 2.18(b). At zero current bias the P configuration of the ND and AF spins is stable, whereas the AP configuration is not. The AP state can exist when either a magnetic field or a bias current producing high STT are applied. When such field or current is decreased to zero the spin configuration of the ND is governed by the exchange coupling with the AF layer, which will force the ND spins to switch into the original parallel state. It therefore remains unclear why the resistance has two stable states at zero current.

2.3.2.1 Effect of the stress anisotropy (papers II, IV)

Our measurements of the STT effect in PCs to single ferromagnetic films confirm the presence of the hysteretic CIMS, originally reported in [64]. An example of a hysteretic IVC for a Cu(20 nm)/Co(5 nm)/Ag(tip) point contact is shown in Fig. 2.19. For this particular PC the MR is about 2% at zero magnetic field and is reduced to 1.8% at 0.25 T. The application of the field also affects the critical switching current. At 0.5 T the hysteresis disappears completely. In order to explain this switching one could apply the model of T. Y. Chen et al. [64], which was reviewed

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8Here P and AP mean two orientations of the ND with respect to the AF layer, and not the P and AP configurations between the ND and the interior spins which actually gives the MR in this model!

9The observed IVC is highly reproducible and returns to the original shape if the field is turned off.
in the previous section. However, there are several reasons why this model is not suitable in our case:

- The model with a vertical nanodomain is suitable for bulk Co rather than thin films. Our thinnest films exhibiting hysteresis are 2 nm thin, which is smaller than the exchange length in Co (∼ 5 nm). This makes a vertical nanodomain highly unlikely.

- We cool down our samples to 4.2 K in zero magnetic field. In order to establish exchange anisotropy with the assumed antiferromagnetic surface layer, a field cooled procedure was required [64, 66], [67, p. 442]. We observe hysteresis in zero-field cooled samples.

- When a metallic PC is created any oxide layer that might have been present at the surface in the contact region is destroyed. Otherwise the contacts are highly ohmic. Thus, the metallic nature of studied contacts is an evidence that there is no oxide at the surface in the PC core region.
• Another consideration is that the proposed exchange-bias anisotropy would be unidirectional, and therefore $R(B)$ and $dV/dI(B)$ hysteretic loops would be expected to shift with respect to $B = 0$ T [67, p. 442]. We, however, routinely observe PCs with a nearly symmetric field dependence of the MR (see e.g Fig. 2.20(a)).

• Surface exchange bias can be diminished by coating the ferromagnetic films with a nonmagnetic metal. Contrary to [64], where some samples were covered with a 2 nm Au layer and showed no hysteresis, we observe hysteretic switching in samples coated with Au and Cu, as shown in Fig. 2.20(b).

• As discussed on p. 30, the micromagnetic origin of the bi-stable nanodomain is hard to couple with the antiferromagnetic exchange anisotropy within the model given by Chen et al. [64].

Figure 2.20: Hysteretic differential resistance for a PC to an uncoated Co film (a) and a Co film coated with 3 nm of Au (b).

The considerations above motivated us to look for alternative explanations. To understand the micromagnetic structure in a nano-sized PC we studied the magnetic field dependence of the resistance at a relatively low current through the PC, in order to minimize the current-induced STT. $dV/dI$ as a function of $B$ at $I = -0.25$ mA ($I_{crit} = -0.75$ mA) is shown in the inset to Fig. 2.21. A distinct characteristic of the field driven hysteretic loop is that the ground state at zero field is a high resistance state. For the field sweep from +200 mT to -200 mT (blue circles) an increase in $dV/dI$ is observed before zero field is reached, starting at about +140 mT.
Figure 2.21: Hysteretic current induced switching in a Cu(100 nm)/Co(5 nm)/Ag (tip) PC having 2.9 Ω resistance measured in zero field. Field induced switching at -0.25 mA biasing current (inset).

Sweeping the field in the opposite direction (red squares) reveals again an abrupt transition to a high-R state before zero field is reached, at a negative field of -64 mT. The question is what kind of magnetic anisotropy can cause an increase in resistance or even an abrupt switching after saturation ($|B| > 1$ T), before $B$ reaches zero?

**Our micromagnetic model vs. experiment**

Here we analyze mechanism for the observed hysteresis effect for PCs to single N/F interfaces, which is *magneto-elastic* in nature. A thin Co film can be subject to a mechanical stress when a metallic tip is brought into contact with the film surface. The resulting magnetoelastic energy $E_{me}$ can be written as [67, p. 270]:

$$E_{me} = \frac{3}{2} \lambda_s \sigma \sin^2 \varphi,$$

(2.22)

where $\lambda_s$ is the linear saturation magnetostriction, $\sigma$ the applied stress, and $\varphi$ the angle between the direction of the stress and the magnetization. We propose that a nanodomain (ND) can be formed when $E_{me}$ in the PC region is high enough to overcome the shape anisotropy of the film. Using Eq. 2.22 we can estimate $E_{me}$ by assuming $\sin^2 \varphi \approx 1$ (for a vertical ND) and $\lambda_s \sim -10^{-5}$ for Co [68, p. 50]. The compressive (negative) stress $\sigma$ is unknown, however the maximum value is the elastic limit for Co, so we can use the Young’s modulus of $-2.1 \times 10^{11}$ Pa [68, p. 25]. These give $E_{me} \approx 3 \cdot 10^6$ J/m$^3$, which is comparable to the demagnetization energy
in Co, $K_d = 1.4 \cdot 10^6$ J/m$^3$.

To test the above hypothesis in detail we used the OOMMF micromagnetic software package [69]. From the numerical micromagnetic analysis we determined the minimal $E_{mc}$ required for a formation of a ND, the ND size, and the estimated MR (a function of the ND-to-film spin angle) as a function of an externally applied magnetic field. For the simulations we chose a 200 nm diameter, 5 nm thick film element, with a 1 nm$^3$ mesh size. Bulk Co saturation magnetization, $M_s = 1.25 \cdot 10^6$ A/m, and exchange stiffness $A = 3 \cdot 10^{-11}$ J/m were used. The stress was taken to fall off symmetrically away from the PC core, with a Gaussian profile of different decay lengths, as shown in Fig. 2.22(a).

Characterized by the value of the anisotropy energies in the center $K_u(r = 0) = K_0$ and at the distance equal to the radius of the PC $K_u(r = r_{pc}) = K_{pc}$ (red vertical line in Fig. 2.22a). The typical value used in the simulations, $r_{pc} = 10$ nm, corresponds to the size of our larger contacts ($R \approx 3 \Omega$). The direction of the anisotropy easy axis was taken perpendicular to the film plane since the easy axis of stress anisotropy in Co is known to be parallel to the direction of the compressive stress.

The simulations results for three stress profiles with $K_{pc} = 0.5 \cdot K_0$, $0.9 \cdot K_0$, and $0.99 \cdot K_0$ are presented in Fig. 2.22(b). $\theta$ is defined as the angle between the magnetization vectors $M_0$ and $M_1$ averaged across the PC region ($r \leq r_{pc}$) and its periphery ($r_{pc} < r \leq 2r_{pc}$), respectively, as shown in Fig. 2.23 inset. Since the current density is high only in the close vicinity of the PC the analysis of the spin distribution is limited to the region of $\approx 2r_{pc}$. For $K_{pc} = 0.5K_0$, $0.9K_0$ and $0.99K_0$, $\theta$ reaches a maximum at energies $0.71K_d$, $1.07K_d$ and $2.5K_d$, respectively. This means that a stable out-of-plane ND corresponds to the magneto-mechanical equilibrium configuration in the system.
For further analysis we choose the medium value with $K_{pc} = 0.9K_0$ (black solid line in Fig. 2.22). At low energies $K_0 \ll K_d$, $\theta$ is zero since the stress anisotropy is much weaker than the shape anisotropy of the film, and no ND can be formed. On the other hand, if $K_0 \gg K_d$ then the size of the ND is much greater than the PC size. Here again, $M_0 \approx M_1$, however the direction is out of the plane. When $K_0 \sim K_d$, the ND size is comparable with the size of the larger contacts ($\geq 20$ nm), $\theta \neq 0$, $M_0 \neq M_1$, and current flows through an effective domain wall producing magnetoresistance.

Our model can be used to obtain the hysteretic MR dependence on the magnetic field. Using $K_0 \approx 1.07K_d = 1.5 \cdot 10^6$ J/m$^3$, $K_{pc} = 0.9K_0$ and choosing $r_{pc} = 10$ nm we compute $\theta(B)$. Since $\text{MR} \sim \sin^2(\theta/2)$ the hysteretic dependence of MR on $B$ is qualitatively similar to $\theta(B)$ shown in Fig. 2.24(a). In this particular case, the maximum $\theta$ is $\approx 20^\circ$. Some contacts we have measured, such as the one shown in Fig. 2.24(b), exhibit hysteretic MR loops in excellent qualitative agreement with the simulated ones (Fig. 2.24(a)).

An important conditions for observing MR is that the current must flow through the DW, which is on the periphery of the PC, and the density of the current should be high enough to perform the current-induced magnetization reversal of the ND. A numerical simulation of the current flow injected through a PC into a Co(5 nm)/Cu(20 nm) bi-layer shows that the current density $j$ in Co is about 10–20% of the current at the
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\[ K_{pc} = 0.9, K_p = 1.5 \times 10^6 \text{ J/m}^3, r_{pc} = 10 \text{ nm} \]

\[ K_p = 3.25 \Omega \]

\[ B \text{ (mT)} \]

\[ \theta (\text{deg}) \]

\[ \rho_{Co}/\rho_{Cu} = 5 \]

\[ d_{PC} = 20 \text{ nm} \]

\[ r \times 10^{-2} (\text{nm}) \]

\[ \rho_{Co} = 5 \]

\[ \rho_{Cu} \]

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Figure 2.24: Simulated $\theta$ versus $B$ applied in the plane of a 5 nm thick Co film (a). Measured MR of a PC with $R = 3.25 \Omega$ (b). Arrows and corresponding color indicate the field sweep directions.

Figure 2.25: A numerical simulation of the current flow in a Co(5 nm)/Cu(20 nm) film, injected through a PC (a). The current density $j$ is shown as a color map. In the limit when the mean free path $l$ is comparable with the PC size, $j$ is a function of the distance $r$ from the PC center (b). Results are shown for $a/l$ equal to 0.1, 0.2, 0.5, 1.0, 2.0 and 10.0, after [70].

N/F interface, as shown in Fig. 2.25(a). In these simulations the ratio between the resistivities of the two metals is assumed\(^\text{10}\) to be $\rho_{Co}/\rho_{Cu} = 5$. From the figure we see that most of the current will traverse the Co layer and enter the Cu buffer layer. Thus, the DWMR is sensed by a small fraction of the longitudinal current flowing in the Co layer nearly parallel to the film. This fact can be used to explain relatively

\(^{10}\) At liquid He temperatures this ratio can be 5-10 for our Co/Cu films.
low MR observed on the experiment, though the actual current density at the PC periphery can be somewhat higher, as was shown by MacDonald and Leavens [70]. If the contact size is smaller or comparable to the elastic mean free path in the metal (Sharvin limit) then the current density increases with distance \( r \) from the PC center and reaches its maximum value at \( r = r_{\text{PC}} \), as shown in Fig. 2.25(b).

Noticeable quantitative differences (Fig. 2.24) between the experimental data and the results of the micromagnetic simulations can be due to the uncertainties in the geometry and material parameters of the PC’s, as well as small but nonzero STT affects. Our micromagnetic model provides a microscopic explanation of the observed effect for larger point contacts (\( \gtrsim 20 \) nm). Hysteresis in smallest contacts (1–10 nm range) may require invoking other mechanisms, since such volume-like nano-domain states on this scale are highly unlikely.

2.3.2.2 Spin valve surface effect (paper III)

A large number of magnetic point contacts we have measured resulted in a large library of magnetotransport data. Processing these data, in particular small PCs which exhibit current- (CDHS) or field-driven (FDHS) hysteretic switching, revealed few different types of the either characteristics.

The FDHS shown in Fig. 2.26(b, d) cannot be explained within the stress anisotropy model presented in section 2.3.2.1. The first distinctive characteristic of this field-induced switching is that the hysteretic loop is qualitatively similar to the typical hysteretic switching for a typical F/N/F spin-valve device (see Fig. 2.8(inset)). Intuitively, this fact suggests that thin ferromagnetic films can be split into two magnetically decoupled (or weakly coupled) layers. We assume that spins at the surface and the interior spins in the ferromagnet can be considered as such layers. An experimental support of this assumption comes from the recent studies of surface and interface magnetism in Co and Fe [71, 72]. It is believed that the difference in magnetic anisotropy and/or exchange coupling is due to the reduced site symmetry and the different atomic coordination at the surface with respect to the bulk.

Another interesting property of the two apparently similar\(^{11}\) contacts in Fig. 2.26 is that the hysteretic \( R(V) \) dependence can be inverted in resistance, switching into the high–\( R \) state (AP) for positive current bias\(^{12}\). Thus, the hysteresis shown in Fig. 2.26(c) is inverted with respect to the typically observed normal hysteresis shown in Fig. 2.26(a), where the AP state is produced by negative current bias. Such inverted hysteresis we will refer to as anomalous.

In point contacts to bulk or thick film ferromagnets it is natural to observe normal CDHS. The reason for that is because there is only one active N/F interface, as shown in Fig. 2.27(a). When the electron current flows from thick F to N (positive polarity), it is spin-polarized and any DW or nanodomain near the N/F point

\(^{11}\)In PC size and film thickness.

\(^{12}\)Here, the negative polarity of \( I \) or \( V \) corresponds to the situation when the electrons flow from N to F at the N/F interface.
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Figure 2.26: Current- (a) and field- (b) driven normal hysteresis for a Cu(30 nm)/Co(5 nm)-Ag(tip) PC. (c, d) same for a similar contact exhibiting anomalous hysteresis.

contact is suppressed if the current, and therefore the STT, is strong enough. In the case of CDHS this corresponds an AP-to-P switching at approximately +2.5 mV in Fig.2.26(a). For the negative polarity of the current the unpolarized electrons flow from N to F. As described in sec. 2.3.1 (paper I), efficient spin-dependent electron scattering on impurities near the N/F interface causes STT acting on the interface spins, which are shown by red arrows in Fig. 2.27(a). The STT may lead to the a reversal of the surface magnetization P-to-AP seen in the data at approximately -11 mV (Fig. 2.26(a)).

The normal hysteresis is observed in our PCs to bulk Co as well as Co thin films (see Fig. 4(a) in paper III). In contrast, the anomalous hysteresis is not observed for bulk Co, but is observed for thin films with thickness of the F film smaller than the PC diameter ($t_F \lesssim d$). Our explanation of this anomalous effect is given by the following model.

The geometry of a N1/F/N2 point contact can be such that the F layer can be located in the bottom (Fig. 2.27(b)) or top (Fig. 2.27(c)) half of the PC core. As illustrated, either interface (N1/F or F/N2) can be in the middle of the PC core (dashed circle) and therefore subject to a higher current density than the other
interface. In these two cases most of the PC core volume will be filled predominantly by N1 or N2 metal. Fortunately, PCS can help to determine which N metal is dominant in the PC core, and correlate that with the two types of the observed CDHS. In Fig. 2.28 the result of such an experiment is presented. For two similar Cu(100 nm)/Co(3 nm)–Au(tip) contacts both CDHS in $dV/dI$ (figures 2.28(a, c)) and $d^2V/dI^2$ spectra (figures 2.28(b, d)) were recorded. The first PC shows normal hysteresis and exhibits a pronounced phonon peak of Au, approximately at 10 mV. In contrast, the second PC shows anomalous hysteresis and, importantly, has a pronounced phonon peak of Cu rather than Au, at 17 mV. The peak separation of about 7 mV between Au and Cu is large enough to distinguish them even in the diffusive regime, where the peaks are smeared out. The choice of the non-magnetic metals N1 and N2 for the electrodes in this experiment was aimed to test our model. The spectroscopic data (d) provide a good explanation. Namely, the PC core with the highest current density is at the bottom (top) of the F/N2 (N1/F) interface, where the electrons flow from Cu (Au) electrode into Co. Thus, transition to the high resistance state occurs at positive (negative) polarity of the current, and the reverse switching into the low resistance state occurs at the neg-
Figure 2.28: Normal hysteresis in $dV/dI(V)$ for a Cu(100 nm)/Co(3 nm)-Au(tip) PC (a). PC spectrum for the same PC showing a pronounced Au transverse phonon peak (c). Anomalous hysteresis in $dV/dI(V)$ for a nominally similar contact (c) and its PC spectrum showing a dominant Cu phonon peak (d). The black arrows in dashed circles in (a, c) schematically indicate the orientation of the surface and interior spins. The red and blue arrows in (a, c) indicate the bias sweep direction.

2.3.2.3 Magnetic vortex states in point contacts

A subset of hysteretic point contacts exhibit CIMS with three stable states in the $R(V)$. These measurements were done on thicker ferromagnetic films ($t_F \gg d_{PC}$). In contrast to the usual, bi-stable hysteresis (e.g. shown in Fig. 2.19, 2.20, 2.21 and 2.26), a third resistance branch appears between the maximum and minimum resistance states. An example of the experimental tristable hysteresis is shown in Fig. 2.29 [65]. The high and low resistances we denote as $R_{AP}$ and $R_P$, respectively, assuming that the AP and P states are realized between the surface spins and the...
Figure 2.29: Measured $dV/dI(V)$ at $T = 4.2$ K and $H = 0$ T for Cu(100 nm)/Co(100 nm)/Cu(3 nm)-Cu-tip PC with total 6 cycles. Each cycle is labeled by color. Only in cycle no. 2 the intermediate resistance state is reached. Black arrows show sweep direction. AP, P, V and “s” represent anti-parallel, parallel, vortex states and beginning of the curve recording respectively. The magnetic transitions between either states are shown in the plot legend.

Bulk. Exact 180° (AP) or 0° (P) alignment of the spins between the two regions is not a necessity. Depending on the PC geometry and the strength of the exchange coupling between the surface and the bulk spins, other angles can provide stable configurations of the spins in the system. The middle resistance state we denote as $R_V$, suggesting that in this state the spins at the surface of the ferromagnet have a vortex micromagnetic configuration.

Following the surface spin-valve interpretation discussed in previous section (2.3.2.2 and paper III), we propose that vortex state can be formed at the surface of F layer, at the top N/F interface (Fig. 2.27(a)), where the current density is the highest. The bottom F/N interface is far from the PC core and plays no role in this case.

A similar $R(V)$ characteristics (Fig. 2.29 vs. 2.30(b)) was found in $F_{\text{ring}}/N/F$ nanopillars [73], where one of the F layers has the shape of a circular ring, as schematically shown in Fig. 2.30(e). Despite the fact that the rings in [73] had lateral dimensions $\varnothing > 100$ nm (an order of magnitude larger than our PCs) remarkable similarity is evident from the $R(H)$ loops for some of our PCs shown in figures 2.31(d–f) and 2.31(b). The middle resistance step in $R(H)$, which has been identified with a vortex state of the nanoring in [73], is apparently present in PC.

The $R(V)$ dependence of the PC shown in Fig. 2.31(a) has three distinctive re-
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Figure 2.30: After T. Yang et al. [73], CIMS loops for ring-shaped nanopillars, schematically shown in (e), with the sizes of (a) 155/60 nm and (b) 215/95 nm. Illustration of the transitions between onion-AP and vortex states (c, d). $M_1$ and $M_2$ (solid lines) are magnetizations of bottom fixed layer and Co ring respectively. The Oersted field and DC current are shown by dotted and dashed lines respectively.

... resistance states. At the start of the sweep (s) the magnetic configuration of the PC is P (black curve 1). At a rather large bias ($\sim -65 \ldots -68\text{mV}$) the magnetic configuration changes to AP and remains in this state until it finally switches into an intermediate resistance V state (green curve 3). The state switches again to AP somewhere between $-20\ldots -40\text{mV}$. In the following half-cycle (blue curve 4) the state remains AP and at about $+15\ldots +20\text{ mV}$ becomes P. Then between $-20\ldots -64\text{ mV}$ the magnetic configuration changes to V. Since in this sweep bias $>-64\text{ mV}$ is not reached, there was not enough current density and thus STT to perform the P-AP transition. Therefore, when the voltage is decreased back to zero the spins remain in the V state (pink curve 6). Using the fifth voltage sweep (cyan curve) one can estimate the current required to form the vortex state (by means of generated Oersted field), which is $\sim 10\text{ mA}$. As detailed below (Fig.2.33), this suggests that the exchange stiffness at the surface must be two orders of magnitude smaller than in bulk Co, if the PC diameter is $\sim 20\text{ nm}$.

**Micromagnetics of vortex nucleation and stability in thin disks**

In our numerical model we study the nucleation as well as stability conditions for vortex-like spin states in thin circular disks. We perform micromagnetic simulations using OOMMF [69]. The equilibrium magnetic configuration is obtained by an energy minimization.

The computational volume is a 1 nm thick disk having the diameter of ei-
ther 20 or 50 nm. The volume has discretized into a three-dimensional mesh of 0.5 × 0.5 × 0.5 nm³ cubic cells. We took the saturation magnetization of $M_s = 1.25 \times 10^6$ A/m and the uniaxial anisotropy energy of $K_u = 0$ J/m³, which are typical for thin Co films. We varied the exchange energy up to $J_{\text{EX}} = 3 \times 10^{-11}$ J/m (bulk Co). Theoretical studies on the vortex stability [74] show strong dependence on the $L/L_{\text{EX}}$ ratio, where $L$ the size of the disk and $L_{\text{EX}} \sim \sqrt{J_{\text{EX}}}$ is the exchange length in the material. In ultra-thin magnetic films with thickness of the order of a few atomic layers, the exchange energy is known to be lower than in the bulk.

To investigate the vortex stability in the PC we use two approaches. In the first one, we determine the vortex nucleation probability $P$ statistically from simulations where in each run the initial magnetization is set randomly from cell to cell. An example of the micromagnetic spin configuration equilibrated to a stable vortex state is shown in Fig. 2.32(a). The results of such computations for disks 20 and 50 nm in diameter are summarized in Fig. 2.32(b). Here we also present the results for the case where a finite out of plane (OP) uniaxial anisotropy was added. As can be seen from the plot this anisotropy does not affect the nucleation probability.

In the second approach we analyze the effect of the current induced Oersted field $H_{\text{Oe}}$ on the nucleation of the vortex. We assume that the current $I$ is uniformly distributed in the PC (magnetic disk) and the field is linearly increasing from zero at the PC center as:

$$H_{\text{Oe}}(r) = \frac{H_0 r}{R_{\text{PC}}}.$$  \hspace{1cm} (2.23)
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(a) vortex 3D view

![3D view of vortex spin distribution](image)

Figure 2.32: 3D view of the vortex spin distribution in a 1 nm thick magnetic disk having a diameter of 50 nm, $M_s = 1.25 \cdot 10^6$ A/m, and $J_{EX} = 1 \cdot 10^{-12}$ J/m (a). This distribution is a result of a micromagnetic simulation where the initial magnetization in each $0.5 \times 0.5 \times 0.5$ nm$^3$ cubic cell is set randomly. The cone-like arrows represent $M$ and the color depicts its normalized $M_z$ component. Vortex nucleation probability $P$ as a function of the exchange energy $J_{EX}$ for 20 and 50 nm diameter PCs (b).

(b) vortex nucleation probability

![Graph showing vortex nucleation probability](image)

Here $H_0 = \mu_0 I/(2\pi R_{PC})$ is the maximum value of the Oersted field at the PC periphery. We start our simulations with the uniform magnetization and the initial $H_{Oe}$ applied field defined by 2.23. After the energy minimization converges to the equilibrium the Oersted field is set to zero and the magnetization state is determined. In Fig. 2.33 the resulting vortex phase diagram is presented for a PC with $R_{PC} = 10$ nm. From this diagram the critical Oersted field required to nucleate a vortex states can be determined and, thus, the critical current can be estimated.

In conclusion, we have observed current- and field-driven hysteretic switching with triple (P, V and AP) resistance levels in PCs to single F films. Qualitatively, both CDHS and FDHS are similar to the STT effects found in spin-valve $F_{mag}/N/F$ structures, where vortex magnetic states have been observed [73, 75]. The typical dimensions of the measured PCs we estimated to be $10-20$ nm, which is very unusual for vortex magnetic states. However, our micromagnetic simulations show that such states are possible in the case of the reduced exchange interaction between the spins at the surface. Our numerical studies also show that creation of the vortex configuration in PCs is stimulated by the Oersted field produced by the high density current.
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\[ \log_{10}(J_{ex}/J_0), \quad J_0 = 3 \cdot 10^{-11} J/m \]

Figure 2.33: Vortex stability diagram, where the Oersted field \( H_{Oe} \) is the minimal field required for the nucleation of a stable vortex in a disk with \( R_{PC} = 10 \) nm. \( H_{Oe}(r = 0) = 0 \) and is maximum at the PC periphery \( H_{Oe}(r = R_{PC}) \), as described in the text. Color diagrams show distribution of the magnetization’s angle in the plane of the disk.

2.3.3 Spin dynamics in magnetic point contacts (paper V)

In the previous sections (2.3.1–2.3.2) we have studied STT effects in point contacts to single ferromagnetic films. We learned that, similar to F/N/F multilayers, single N/F interface can exhibit magneto-transport properties such as resistance excitations and even hysteretic magnetization reversal caused by nominally unpolarized currents. Our experimental confirmation (paper I) of the predicted \([61, 63]\) STT mechanism for single N/F interface helped us to understand the nature of the the magnetic excitations (paper II) and the origin of the current-induced hysteretic switching (papers III, IV and \([65]\)).

So far, the experimental studies we have presented were limited to static measurements. By measuring such dependencies as \( dV(I, H)/dI \) (and \( R(I, H) \)) shown in Fig. 2.34(a) and \( d^2V(I, H)/dI^2 \) we can characterize N/F/N point contacts by the current flow regime, presence of magnetic excitations, and some additional details on the magnetization reversal if the magnetic PC exhibits hysteretic switching. By looking at the progress of the theoretical and experimental research on STT in N/F multilayers or F/N/F spin-valve devices, we see that these studies were accompanied by investigations of STT dynamics. In the phenomenological LLG equation 2.19 the STT, defined by Eq. 2.20, is an additional term which depends on the degree of spin polarization \( P \) and current density \( J \). As discussed in section
2.3. RESULTS AND DISCUSSIONS

Figure 2.34: Relative differential resistance for Co(100 nm)/Cu(tip) PC with $R_0 = 7.2 \Omega$ with $H_\perp = 4$ T applied field (a). The schematics of the experiment is shown in inset of (a). Excitation peak in $dV/dI$ stimulated by 2 GHz irradiation of different power $P = 0, 2.4$ and $3.6 \mu W$ (curves 1–3 respectively) for Co(100 nm)/Cu(tip) PC with $R_0 = 5 \Omega$ in $H_\perp = 2.47$ T field (b). Curves are shifted vertically by $1 \Omega$ for clarity. The amplitude of the induced peak as function of RF power is shown in the inset of (b).

2.1.4.2, the STT is responsible for the magnetization excitations in ferromagnetic free layers induced by the polarized current in the spin-valve. In fact, when $J$ is approximately equal to a certain critical value $J_c$, the STT causes a steady-state magnetization precession, schematically illustrated in Fig. 2.7. This precession can be indirectly observed as peaks in static resistance measurements. The direct evidences of the STT-induced magnetization dynamics in F/N/F trilayers has been obtained in [76]. For single F layers such direct measurements remain a challenging experiment. We point out that the first evidence that the magnetic excitations in multilayers are oscillatory in nature was obtained using quasi-dynamic measurements in [77]. Similar to this work, here we demonstrate STT-induced magnetization dynamics in single F layer.

The schematic of the experiment, shown in Fig. 2.34(a)(inset), is not very different from the previous experiments with point contacts discussed in the previous sections. The same four point measurement technique is used to inject high density current into F from a N tip and to sensitively measure the $dV/dI$ characteristics. The novelty with this experiment is that, a microwave radiation is supplied via a coaxial cable to the tip. The DC and HF lines are decoupled by using a capacitor and low ohmic inductors, as shown in the figure. In this case the PC needle plays the role of an efficient antenna at GHz frequencies, supplying a HF AC cur-
rent to the PC. The maximum RF power in the PC region is estimated to be a few tens of µW.

The main idea behind this RF irradiation experiment is to probe the HF frequencies of the current-induced magnons. One expects expected to observe a resonant amplification of the current-induced magnetization precession (peaks in $dV/dI$) when the frequency ($f_{RF}$) of the externally applied RF coincides with the magnon precession frequency ($f_m$) - a function of the bias current through the PC, the local magnetic anisotropy and externally applied field. When the resonant condition is met a transition from predominantly stochastic oscillations to a stationary precession occurs. The energy of the spin sub-system in the PC core increases with respect to the spins in the bulk (film). This may lead to an abrupt change in the DWMR seen as a peak in $dV/dI$. As an example, a stimulation of magnetic excitation by RF is shown in Fig. 2.35(a), where the peak appeared first when the PC was irradiated by RF of 3 GHz power of 4 dBm\textsuperscript{13}. An increase of the frequency caused a shift of the peak to a higher bias. Importantly, a shift towards a higher bias is observed by increasing the external field (Fig. 2.35(b)).

In addition to the resonant absorption of RF discussed above, we also observe off-resonance AC rectification. An example of this effect is shown in Fig 2.35(c) where initially sharp peak in $dV/dI$ (curve 1) is suppressed, broadened and then even split in two with an increase of the RF. Such behavior can be studied theoretically by using the model given of Refs. 17 and 18 in the manuscript (paper V). Since the total voltage over the PC is $V(t) = V_0 + V_1 \cos \omega_{RF}t$, where $V_0$ is the DC bias and $V_1$ is the amplitude of the RF with frequency $\omega_{RF}$, the time averaged IVC characteristics is:

$$I(V) = \frac{\omega_{RF}}{\pi} \int_0^{\pi/\omega_{RF}} I_0(V_0 + V_1 \cos \omega_{RF}t)dt.$$  \hspace{1cm} (2.24)

Here $I_0(V_0)$ is unperturbed IVC, which can be obtained by integrating the experimentally measured $dV/dI$ without RF applied (curve 1 in Fig. 2.35(c)). Eq. 2.24 requires that the RF frequency is low compared with the inverse of the characteristic electron relaxation time producing nonlinearity. In other words, energy of the photons must be lower than the width of the nonlinearity in the IVC. In our experiments this requirement indeed holds, the energy of the RF photons is of the order $\hbar \omega_{RF} \sim 10^{-2}$ meV, whereas the half-width of the peak is about $\sim 2$ meV. Now, having an unperturbed $I_0(V_0)$ dependence and using Eq. 2.24 we can compute $T(V)$ – the IVC for different values of RF power (voltage amplitudes). The result is shown in Fig. 2.35(d) (curves 2–5), which is in prefect agreement with the experimental curves in Fig. 2.35(c).

In some contacts the amplitude of the peak was monotonously rising as RF

\textsuperscript{13}Measured at the RF generator.
2.3. RESULTS AND DISCUSSIONS

Figure 2.35: Differential resistance for a PC with \( R_0 = 6.7 \) \( \Omega \) (a) and 18 \( \Omega \) (b) irradiated by RF of given frequencies, with power \( P_{RF} = 0 \) and 4 dBm. \( \frac{dV}{dI}(V) \) for a Co(100 nm)/Cu(tip) PC subjected to a 2 GHz RF irradiation with power of \( P = 0, 12, 24, 36 \) and 48 \( \mu \)W marked as curves 1–5, respectively (c). Calculated dependence according to Eq. 2.24 for \( V_1 = 0.5, 1.5 \) and 2 \( \mu \)V, marked as curves 2–5, respectively (d). Curve 1 in (c) is the same as curve 1 in (d). Peak position as a function of the external field (inset).

Power increased, as shown in Fig. 2.34(b). Such behavior is expected for a transition from a low to a high angle precession with increasing RF power. As can be seen from Fig. 2.34(b)(inset), for low RF power the amplitude of the induced peak \( A_{AC} \) depends on the power linearly. This is in agreement with recent measurements on spin-transfer-driven ferromagnetic resonance in spin-valve devices [78], where peaks have been well approximated by the Lorenzian shape\(^{14}\) (Eq. 3 in [78]):

\[
A_{C} \propto \frac{P_{RF}^2/\Delta_0}{1 + [(f_{RF} - f_m)/\Delta_0]^2},
\]

\(^{14}\)Our experimental \( A_{C}(f_{RF}) \) dependence is given in paper V. (Fig. 4(b)).
where $\Delta_0$ is the linewidth of the peak. Since $P_{RF} \propto I_{RF}^2$, we can rewrite 2.25:

$$A_C \propto \frac{P_{RF}/\Delta_0}{1 + [(f_{RF} - f_m)/\Delta_0]^2}. \quad (2.26)$$

I. e., the peak amplitude scales linearly with the RF power when power is sufficiently low.

Higher RF power results in a saturation of the precession angle. Driven by the STT and the RF, the surface spins in the PC core form an effective precession angle with spins outside the core, producing DWMR. This DWMR is measured as a pronounced peak in $dV/dI$. 
Chapter 3

Spin-Flop Dynamics in Toggle MRAM (paper VI)

*Toggle magnetic random access memory* (TMRAM) has recently attracted a great deal of attention from many research groups. Invented by L. Savchenko, TMRAM has two antiferromagnetically coupled free magnetic layers separated by a non-magnetic spacer. It has been shown [79, 80], that such a design makes the layers insensitive to the disturbing half-select fields from either word or bit lines, and the switching can be performed when a specific *toggle* sequence of the word and bit line fields is applied. This resolves the problem of unwanted, variations in switching fields. Furthermore, theoretical single-domain modeling [81–83] predicted enhanced activation energy when a half-select field is applied, and recent experimental results confirmed this [84]. Other theoretical [85–87] as well as experimental [88] studies were concentrated on quasi-static TMRAM properties, such as dependence of the toggle critical switching curve (CSC) or the toggle criteria on free layers geometry, thickness imbalance, intrinsic anisotropy, etc.

The magnetization dynamics of a *single* ferromagnetic nanoparticle induced by external fields [53, pp. 84–103], [89, 90] and spin transfer torques [57, 58, 91–93] is well understood. It is governed by the LLG equation (2.19), which describes a spin precession with the resonance frequency determined by the effective anisotropy field \( H_{\text{eff}} \) of the particle.

The dynamics of *two* ferromagnets, coupled either by exchange or dipolar forces or both, is however practically unexplored. In this two-particle case, one should expect collective precessional modes to replace the single particle resonances when the inter-particle coupling strength is comparable to the single particle magnetic anisotropy energy.

In this chapter, we report real time measurements of the tunneling magnetoresistance of spin-flop tunnel junctions in response to 100 ps range excitations. We show that the two coupled nanomagnets, when dynamically excited by a magnetic field impulse, resonate in a collective fashion with the two main eigen-modes cor-
CHAPTER 3. SPIN-FLOP DYNAMICS IN TOGGLE MRAM (PAPER VI)

responding to the collective azimuthal bodily rotation and the azimuthal bending of the two macrospins. The presence of two precessional eigen-modes in the spin-flop system is supported by a theory of its LLG dynamics in the macrospin approximation. We also present results from our micromagnetic numerical studies, showing that the precessional spectra are rather complex and have transitional features which single-domain model cannot explain. The results of high speed measurements agree well with the micromagnetic simulations, and a good agreement with the single-domain model is obtained for higher values of the magnetic field, when the moments of the coupled layers are in the scissored magnetic state.

3.1 Theory

In this section we briefly review the basic idea behind toggle switching and provide theoretical details required to understand how the toggle cell really works. Hereafter toggle cell refers to two coupled ferromagnetic layers rather than a complete memory cell. These two layers are the basic elements of the TMRAM. Most of the magnetic properties of interest for memory can be well described by a single-domain model (SDM). A good overview of the SDM for TMRAM has been recently published by our collaborator Daniel C. Worledge [83].

Since our studies are focused on the fast magnetization dynamics, we also discuss below the analytic as well as numerical results obtained from SDM for the time dependence of the magnetic moments in the system. Later, these results will be compared with our micromagnetic simulations and high-speed measurements.

3.1.1 Toggle switching in spin-flop tunnel junctions

In any type of magnetic storage device a single bit of information, binary «0» or «1», is stored as the magnetization direction aligned P (say for «0») or AP («1») with respect to some reference direction. The physical implementation of such memory bit can be a magnetic domain as in hard drives or magnetic tapes, or a uniform spin state of a nanoparticle in spin-valve devices. On the one hand, magnetic nanoparticles have to be stable against thermal fluctuations and parasitic magnetic fields, which is achieved by selecting a proper magnetic material with suitably high shape anisotropy. On the other hand, the shape anisotropy should not be too high because otherwise the magnetic field required to perform the write operation (magnetization reversal in the particle) is too high. Thus, the choice of the magnetic material and the geometry, specifically the aspect ratio (AR) for elliptical particles, is based on the compromise between the requirement for stability and the need for low switching fields.

The basic MRAM cell consists of a several key elements: a free F layer having elliptical shape, which can be switched in either direction along the length; a non-magnetic insulating spacer separating the free layer from the reference (fixed) layer. This F/I/F free layer is a tunnel junction with TMR ≈ 20% (for AlOₓ spacers). For minimizing the fringing fields acting on the free layer and for a better
3.1. THEORY

thermal stability, the fixed layer is usually replaced by two AF coupled layers, as shown in Fig. 3.1(a)(left). The magnetization reversal of the free layer can be per-

![Diagram](image-url)

Figure 3.1: Comparison of Stoner-Wohlfarth and toggle MRAM (a). The main structural difference is the additional free layer in the toggle design. Word and bit line alignment with respect to TJ easy axes in the toggle design (b) (top). Toggle write BL, WL pulse sequence (b) (bottom).

formed by an externally applied field, which is a vector sum of the bit line (BL) \( H_{BL} \) and word line (WL) \( H_{WL} \) fields, allowing to select a specific memory cell, as shown in figure 3.1(b) top panel. Each BL (WL) is a conductor strip and is located beneath (above) the MTJs. When current is sent through BL/WL a magnetic field is created in the plane of the free layer at the BL-WL cross. This field is directed along the easy axis (EA) of the ellipse. By simultaneously applying current in BL and WL one can switch the direction of the M in the free layer. The reversal of the free layer can be described by the Stoner-Wohlfarth (SW) model. For switching it is sufficient to cross the SW astroid boundary shown, in Fig. 3.2(a).

The main disadvantage of such switching is that the other, not selected cells on the same BL and WL experience so called half-select. The activation energy \( E_a \), the energy barrier preventing the free layer from the thermally activated switching, decreases drastically at half-select (Fig. 3.2(c)). As a result, for half-selected bits a weak thermal fluctuation can lead to a switching error.

The unwanted effect of the half-select described above is minimized in toggle MRAM. In this design, a non-magnetic spacer and then second free layer are added, as schematically shown in Fig. 3.1(a)(right). The critical switching curve changes completely from the SW astroid to the L-shaped toggle boundary.
Figure 3.2: Illustration of the Stoner-Wohlfarth (a) and toggle (b) switching criteria. Switching occurs when magnetic field crosses boundaries of astroid (a) or L-shaped curve (rectangular field excursion) (b). When field is applied along half-select direction (either $H_{BL}$ or $H_{WL}$), activation energy $E_a$ decreases for SW particle (c) and increases in the case of TMRAM (d). After D. C. Worledge [83].

(Fig. 3.2(b)). Now, half-selected bits can tolerate fields much greater than the typical toggle field. Moreover, the activation energy increases with field, as shown in Fig. 3.2(d), making the half-selected cells more rather than less thermally stable.

The magnetization reversal in TMRAM is performed differently from that in SW MRAM. Due to the dipole interaction, the two free layers prefer to stay in the AP configuration at zero field. When a field is applied along the EA, as shown in Fig. 3.3(a) by black dashed arrow, both magnetic moments remain AP until a critical field, $H_{sf}$, is reached. At this field a spin-flop (SF) occurs and the moments discontinuously jump from AP into a scissored state (SS). Which moment rotates clockwise and which counterclockwise is undefined (black arrows in the same figure). Thus, application of the $H_{EA}$ does not change the state of the bit in a controllable way. Upon a further increase of the EA field the angle between the two moments decreases monotonically to zero at saturation field, $H_{sat}$. In order to perform a controllable magnetization reversal of two coupled layers, a rectangular field excursion must be used, as shown in Fig. 3.3(b) by the green line. When such a field excursion is applied, the two moments rotates gradually, toggling the state of the bit. This rectangular field excursion is achieved by applying two pulses...
3.1. THEORY

Figure 3.3: Illustration of a spin-flop transition when the field is increased gradually along the EA (black dashed line) (a). Black solid arrows denote moments in the “undefined” configuration (either of the two scissored states). Blue and red arrows denote well defined AP and SS. The BL-1st or WL-1st half-toggle field excursions are shown by brown and green dashed lines, respectively. A rectangular (toggle) field excursion switches the bit in a controllable way (b). During the excursion moments rotate gradually and the final AP state is the opposite of the initial AP state.

A single-domain model was developed to describe the magnetostatic properties of toggle cells, such as $H_{sf}$, $H_{xsat}$, $ysat$, etc. Here the total energy can be written as [81]:

$$
e(\theta_1, \theta_2) = -h_x(z \cos \theta_1 + \cos \theta_2) - h_y(z \sin \theta_1 + \sin \theta_2) +$$

$$+(n_x - jz) \cos \theta_1 \cos \theta_2 + (n_y - jz) \sin \theta_1 \sin \theta_2 +$$

$$+\frac{z}{2} (n_y - n_x + h_i) \sin^2 \theta_1 + \frac{1}{2z} (n_y - n_x + h_i) \sin^2 \theta_2,$$

(3.1)

where $e = Eb/\pi^2 M_s^2 ab t_1 t_2$, $h_{x,y} = H_{x,y} b/4\pi M_s t_1$, $j = J b/4\pi M_s^2 t_1^2$, $z = t_1/t_2$, $E$ is the energy, $\theta_{1,2}$ are the angles of the two moments with respect to $x$ (EA) direction shown in Fig 3.4(a), $t_{1,2}$ are the thicknesses, and $a$ and $b$ are the length and width of the ellipse in $x$ and $y$ respectively, $M_s$ is the magnetization saturation, $J$ is the exchange coupling constant between the layers, and $H_{x,y}$ are the applied fields in $x$ and $y$. Energy $e$ and fields $h_{x,y}$ are given in reduced units. The expressions for demagnetizing factors $n_{x,y}$ are given in [82] with a relatively weak dependence on AR. Our typical samples have AR=1.2, which corresponds to $n_x \approx 0.62$ and $n_y \approx 0.82.$
For simplicity one can assume that the two layers are identical, and that there is no thickness imbalance, \( z = 1 \). Another important assumption is that the dipole field outside one layer, acting on the other, is equal to the demagnetization field inside the first layer. The latter assumption is justified in the case when the spacer between the two layers is thin. As described in [81], the expression for energy (3.1) can be used to obtain the SF critical field, at which, the system loses stability and spin-flops into the SS from the AP states \( e(0, \pi) \) or \( e(\pi, 0) \). Similarly, the saturation field \( H_{\text{sat}} \) can be derived from the energy for the P state \( e(0, 0) \). The final
expressions for the main critical fields, $H_{sf}$, $H_r$ and $H_{sat,y sat}$ are [83]:

$$H_{sf} = \sqrt{H_i \left( 8\pi M_s n_s t_B - \frac{2J}{M_s t} + H_i \right)}, \quad (3.2)$$

$$H_r = \left( 8\pi M_s n_s t_B - \frac{2J}{M_s t} - H_i \right) \sqrt{\frac{H_i}{8\pi M_s n_y t_B - \frac{2J}{M_s t} + H_i}}, \quad (3.3)$$

$$H_{sat} = 8\pi M_s n_s t_B - \frac{2J}{M_s t} - H_i, \quad (3.4)$$

$$H_{ysat} = 8\pi M_s n_y t_B - \frac{2J}{M_s t} + H_i, \quad (3.5)$$

where $H_r$ is the so called return field. When an applied EA field is decreased from some value (assuming saturation or at least $H_{EA} > H_{sf}$) to zero, the moments do not spin-flop from SS back into AP at $H_{sf}$ because of the hysteresis for $H_r < H_{EA} < H_{sf}$, where energy (3.1) has two minima\(^1\): one for the AP and the other for the SS state. $H_r$ however does not define the toggle criteria, $H_{sf}$ does. In order to toggle the field excursion box shown in Fig. 3.3(b) must contain the $H_{sf}$ point. Further analysis of the toggle criteria [85–87] shows that there is continuum of critical points in the $H_{BL}$, $W_L$ space, where $H_{sf}$ are defined by expressions 3.2 and 3.3 for the easy axis only. If the field is applied off the EA then one has to calculate the minima of (3.1) numerically. The result of such a calculation is shown in Fig. 3.4(b) where a critical switching curve (CSC) has the shape of an astroid. As was mentioned before, inside the CSC both AP and SS states can exist and, as expected, below the return point only, the AP state is stable while above the SF only the SS is stable. Performing a field excursion such that it crosses the boundaries of the CSC leads to discontinuous transitions AP $\leftrightarrow$ SS. These transitions occur: clockwise (CW) or counterclockwise (CCW) rotations of the moments, as shown in the figure. A continuous state transition is possible outside of CSC, when a rectangular field excursion is applied.

The size and position of the CSC is determined by the ratio between the dipole and demagnetizing fields as well as the physical dimensions of the layers. The astroid shape of the CSC shown in Fig. 3.4(b) is for the case when both layers are equal in thickness. If the thickness imbalance is present, then the CSC changes into a heart-shaped figure [86]. The box field excursion is the most reliable way to toggle spin-flop TJs, though there exist other types of switching, such as L- and U-shaped field excursions which can cross the CSC. These types of switching and the corresponding toggle conditions have been studied in [87].

### 3.1.2 Spin-flop dynamics: single domain model

The SDM extended to treat spin dynamics predicts two resonant modes, which are collective modes of the coupled magnetic moments. The first, *acoustical* mode

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\(^1\)There are two additional degenerate minima due to the symmetry of the system.
is associated with the restoring force coming from the intrinsic anisotropy $H_i$ acting on both layers simultaneously (Fig. 3.5). Since the typical $H_i$ values are low

$\Delta \theta = \text{const}$

$\Delta \theta \neq \text{const}$

Figure 3.5: Illustration of the acoustical (left) and optical (right) spin precessional modes. The acoustical mode corresponds to an oscillatory motion with frequency $f_a$, where the magnetic moments rotate in phase and the relative azimuthal angle $\Delta \theta$ between moments is constant. The optical mode corresponds to the moments oscillating with frequency $f_o$, out of phase with $\Delta \theta \neq \text{const}$.

in TMRAM free layers, the resulting acoustical microwave oscillations have relatively low frequencies. The second, optical mode is due to the strong demagnetizing force in the system, resulting in a relatively high frequency of this mode. The AP eigen frequencies are [95]:

\[
H_{EA} = 0 \begin{cases} 
  f_a \approx \frac{\gamma}{2\pi} \sqrt{4\pi M_s H_i}, \\
  f_o \approx \frac{\gamma}{2\pi} \sqrt{4\pi M_s H_{ysat}},
\end{cases}
\]

where $\gamma = 175$ GHz/Oe, $H_{EA}$ is an external field along EA, $H_{i(x,y)\text{sat}}$ are saturation fields defined by Eqs. 3.4 and 3.5.

In the SS, the frequencies of two eigen modes are intertwined and both depend on the demagnetizing force in the system [95]:

\[
H_{SF} < H_{EA} < H_{x\text{sat}} \begin{cases} 
  f_a \approx \frac{\gamma}{2\pi} \sqrt{4\pi M_s \left[ \frac{H_{ysat} H_{EA}^2}{H_{x\text{sat}}^2} - H_i \right]}, \\
  f_o \approx \frac{\gamma}{2\pi} \sqrt{4\pi M_s \left[ H_{x\text{sat}} - \frac{H_{EA}^2}{H_{x\text{sat}}} \right]}.
\end{cases}
\]

Expressions (3.6) are valid for the AP states where $H_{EA} = 0$, whereas the (3.7) was derived for the SS and non-zero field. These formulas are derived for the case when the moments are subject to a small field perturbation from their initial equilibrium state. The $f_{a,o}(H_{EA})$ dependence is plotted in Fig. 3.6 for 350 nm wide,
magnetic particles of 1.2 in aspect ratio, 10 Oe intrinsic anisotropy, and zero inter-layer exchange coupling. As seen from the figure, the acoustical mode (blue curve) increases with field whereas the optical branch (red curve) decreases to zero.

Expressions for the eigen frequencies (3.7) were further verified by numerical simulations within the SDM: we solved the LLG equation (2.19 but without STT term) for $M_{1,2}(t)$, as a response to a field impulse. Angles $\theta_{1,2}(t)$ between the moments $M_{1,2}(t)$ and the EA direction were computed. Then, from the Fourier transform of $\theta_{1,2}(t)$ we obtained $f_{a,o}(H_{EA})$ and compared the results with the analytical formulas. An excellent agreement was obtained [95].

3.2 Experimental techniques

3.2.1 MRAM Samples

The samples used in this work were fabricated by methods similar to those described in [96] and had similar multilayer structure and magnetic properties as those described in [88]. The data presented in this chapter are for elliptical tunnel junctions with lateral dimensions ranging between $350 \times 420$ and $400 \times 490$ nm. The Py top and bottom free layers are 5 nm thick, separated by a 1 nm

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3See also sec. 3.3.1.
TaN spacer. The AlO$_x$ tunnel barrier separates the bottom free layer and the top pinned (fixed) layer. In order to reduce the impedance mismatch between the high resistivity TJ’s and the 50 $\Omega$ high frequency (HF) measurement electronics, the AlO$_x$ tunnel barrier was made thinner compared to the samples used in [88]. The typical resistance of the spin-flop TJ’s used in this work was $R_{TJ} \simeq 2$ k$\Omega$ and TMR = 15 – 20%.

Throughout this chapter the parallel (antiparallel) state of the spin-flop free layers, and those of the bottom free and top fixed layers of the TJ are abbreviated as P (AP) and P$_{TJ}$ (AP$_{TJ}$), respectively.

3.2.2 Measurement technique

3.2.2.1 Electrical measurements

Both quasistatic and high speed measurements were performed at room temperature. The on-chip integrated write and read lines (RL) were contacted using surface probes with DC–40 GHz bandwidth. 100 ps range field excitations (impulses) were applied to the TJ’s through the word lines (WL), designed to have the resistance of approximately 50 $\Omega$ for impedance matching purposes. The bandwidth of the driving circuit was estimated to be higher than 6 GHz, the highest frequency measured in this study. A 5 GS/s arbitrary waveform generator Tektronix® AWG7052 was used to produce WL impulse signals with the rise time of 100 ps. For quasistatic readout the TJ’s were current biased with 50 – 100 $\mu$A and the voltage across the junctions was measured using a multimeter for various external static field configurations as shown by the circuit schematics in Fig. 3.7.

For reading out the dynamic contribution to the TJ voltage a bias tee was used to separate the DC from the microwave signal, which was then amplified by +64 dB and recorded on a 40 GS/s Agilent® DSO80604B real time oscilloscope with a 6 GHz bandwidth. A 5 ns real time traces were averaged 1024–4096 times to improve the signal to noise ratio. A toroidal$^3$ bi-axial magnet [88] oriented such that $H_{magnet} \parallel H_{BL,WL}$ was used to produce various static parallel, antiparallel, and scissors configurations of the spin-flop trilayer.

Beside the wideband noise in the system, the main experimental obstacle in the measurements was an inductive pickup across the BL-RL circuit used for readout. The WL impulses and the associated RF fields induced a voltage in the readout circuit because of: the close proximity of the WL and BL-RL. The actual signal of the magnetic origin from the SF TJ was two orders of magnitude lower than the inductive pickup. Fortunately, during the first 5–10 ns right after the WL impulse the pickup was a weak function of time. In order to subtract this voltage from the total signal we used the following procedure. The total averaged TJ voltage response to a WL field (impulse) can be written as:

$$
\langle V_{tot}(t, I_{TJ}, H_{ext}, H_{imp}^{WL}) \rangle = \langle S_{TJ}(t, I_{TJ}, H_{ext}, H_{imp}^{WL}) \rangle + \langle S_{p}(t, H_{imp}^{WL}) \rangle,
$$

3This magnet was made by Erik Lindgren, as part of his MS thesis [94, Ch. 13.0.4].
3.2. EXPERIMENTAL TECHNIQUES

Figure 3.7: Schematic of the electrical circuit used in high speed measurements of the magnetization dynamics in spin-flop tunnel junctions.

where \( t \) is time, \( H_{\text{ext}} \) externally applied field, \( H_{\text{WL}}^{\text{imp}} \) WL field excitation, \( I_{\text{TJ}} \) DC bias through the TJ, \( \langle S_{\text{TJ}}(t, H_{\text{ext}}, H_{\text{WL}}^{\text{imp}}) \rangle \) the signal of the magnetic origin coming from the oscillations of the free layers, and \( \langle S_{p}(t, H_{\text{WL}}^{\text{imp}}) \rangle \) is the inductive background. \( \langle \ldots \rangle \) means averaging, which reduces the noise. Each measurement was done twice, once at a positive TJ bias, \( +I_{\text{TJ}} \), and once at \(-I_{\text{TJ}}\):

\[
\langle V_{\text{tot}}(t, +I_{\text{TJ}}, \ldots) \rangle = \langle S_{\text{TJ}}(t, +I_{\text{TJ}}, \ldots) \rangle + \langle S_{p}(t, H_{\text{WL}}^{\text{imp}}) \rangle = \langle S_{\text{TJ}}(t, +I_{\text{TJ}}, \ldots) \rangle + \langle S_{p}(t, H_{\text{WL}}^{\text{imp}}) \rangle, \tag{3.9}
\]

\[
\langle V_{\text{tot}}(t, -I_{\text{TJ}}, \ldots) \rangle = \langle S_{\text{TJ}}(t, -I_{\text{TJ}}, \ldots) \rangle + \langle S_{p}(t, H_{\text{WL}}^{\text{imp}}) \rangle = -\langle S_{\text{TJ}}(t, -I_{\text{TJ}}, \ldots) \rangle + \langle S_{p}(t, H_{\text{WL}}^{\text{imp}}) \rangle. \tag{3.10}
\]

Now, a subtraction of 3.10 from 3.9 leads to:

\[
\langle \Delta V_{\text{tot}}(t, I_{\text{TJ}}, \ldots) \rangle = 2\langle S_{\text{TJ}}(t, I_{\text{TJ}}, \ldots) \rangle. \tag{3.11}
\]

In principle, one can subtract \( \langle V_{\text{tot}}(t, +I_{\text{TJ}}, \ldots) \rangle \) and \( \langle V_{\text{tot}}(t, I_{\text{TJ}} = 0, \ldots) \rangle \), however the advantage of 3.11 is that we get an amplification of the signal by a factor of...
two\(^4\). Our main assumption is that the background \(\langle S_p(t, H_{\text{WL}}^{\text{imp}}) \rangle\) does not depend on biasing current. Naturally, two measurements with \(\pm I_{\text{TJ}}\) must be performed at the same external magnetic field and the amplitude of the WL impulse should be small enough not to cause a magnetization reversal or change the state of the spin-flop TJ.

3.2.2.2 Data post-processing

The \(\langle \Delta V_{\text{tot}}(t, H_j) \rangle\) dependence, recorded when magnetic field, \(H_{\text{EA}}^{\text{min}} < H_j < H_{\text{EA}}^{\text{max}}\), \(j \in [1, N]\) is applied, was used to compute spectrograms. For a better visual recognition of the peaks in the spectrum a Hann window \([97, \text{p. 553}]\) was applied for each \(\langle \Delta V_{\text{tot}}(t, H_j) \rangle\) data set prior to the computation of the power spectral density (PSD) \(P_{H_j}(f)\). The PSD was computed by using a standard fast Fourier transform. The amplitude of the acoustical and optical peaks in the PSD can vary significantly. Therefore, some of the PSD\(_{H_j}(f)\) figures in this chapter are plotted in logarithmic scale, to bring out as much details in spectral details as possible.

In order to minimize possible inconsistencies in data processing, the same procedure and code were used to compute the spectrograms from the micromagnetic simulations data. Instead of a voltage time dependence, in the simulations we get the magnetizations components \(M_{1,2}^{x,y,z}(t, H_j)\) as functions of time, for each layer. Since we want to compare these results with the experimental data we have to take into account the effect of the TJ conductance. Since the TJ resistance is well approximated by 2.17, we can write the voltage (resistance) across TJ as:

\[
V(\theta) = \frac{1}{1 + P^2 \cos \theta} \text{ (a. u.),}
\]

where polarization \(P = 0.3\), \(\theta(t, H_j)\) is the angle between either layer’s magnetization vector \(M_{1,2}^{x,y,z}\) and \(x\) or \(EA\), as shown in Fig.3.4(a), \(\theta(t, H_j)\) is first calculated from the magnetization components for each layer. Then, from 3.12, we can compute the PSD in the same way as for the experimental data.

A simple estimate using 3.12 of the voltage response \(\delta V\) due to a small magnetization rotation of the layer, \(\delta \theta\), shows that in the \(P_{\text{TJ}}\) (say, \(\theta \approx 0\)) or \(AP_{\text{TJ}}\) (\(\theta \approx \pi\)) states \(\delta V\) can be \(10^1 - 10^2\) times smaller than that for intermediate values of \(\theta\). This results in poor voltage sensitivity for the \(P_{\text{TJ}}\) and \(AP_{\text{TJ}}\) states, and fairly good RF voltage response for the SS. Our experimental results confirm this consideration (see sec. 3.3.2) – most of the HF data presented herein were measured in the SS configurations, only weak traces of spin oscillations were detected in the AP state.

\(^4\)Though our measurements with zero bias gave essentially the same result.
3.3 Results and discussions

3.3.1 Spin-flop dynamics: micromagnetic simulations

The main idea behind the micromagnetic simulations ($\mu$MS) we have performed was to extend the theoretical study of the SF dynamics of two coupled ferromagnets beyond the single-domain model. $\mu$MS explained some spectral features found on the experiment, which could not be understood within the SDM.

We use the micromagnetic framework OOMMF [98] developed by NIST. Since we are interested in time evolution of the magnetic system we configure OOMMF to solve the LLG equation for $M(t)$, with material properties typical for Py: saturation magnetization $M_s = 8.4 \cdot 10^5$ A/m, exchange stiffness $A = 1.3 \cdot 10^{-11}$ J/m, and damping constant $\alpha = 0.025$. The intrinsic anisotropy $H_i$ varied in our simulations from a few Oe up to a few tens of Oe. We define the geometric volume as two 5 nm thick elliptical disks separated by 5 nm thick spacer. The entire simulation volume is discretized by a $5 \times 5 \times 5$ nm$^3$ rectangular mesh. The disks simulated had the same lateral dimensions as those measured on experiment. To test our micromagnetic model and compare it with SDM we also simulated $120 \times 100$ and $60 \times 50$ nm disks.

The initial magnetization of the system can be arbitrary. We set $M_{\text{init}}^1$ and $M_{\text{init}}^2$ at some angle to each other (neither 0 nor $\pi$). The first 6 ns of the simulation time is typically more than enough for the system to equilibrate to a stable state. After that we apply a field excitation, which is a Gaussian shaped impulse shown in Fig. 3.8(a)(top). Since we are interested to study how the SF dynamics depends on the magnetic field, the total magnetic field is a sum of the external DC field components $H_{BL}$, $H_{WL}$ and the impulse field applied along the WL. In the example in Fig. 3.8(a)(top) $H_{EA} = 80$ Oe and the amplitude of the impulse is 20 Oe. The impulse field shown in the figure is negative, however applying the impulse in the positive direction does not change the dynamics of the system. The impulse half-width was chosen to be 200 ps, which is the same as in our experiments.

As explained in sec. 3.2.2.2, instead of the magnetization components $M_{x,y,z}$, we have to compute a quantity, with the help of which we can characterize the magnetization oscillations and make a direct comparison with the experiment. This quantity is the resistance (voltage) of the TJ, computed according to Eq. 3.12. On the experiment, either of the magnetic layers can face the TJ, thus we analyze spectrum for each layer.

The spectrograms obtained from four simulations for $60 \times 50$, $120 \times 100$, $450 \times 375$ and $560 \times 400$ nm elliptical bi-layers, excited by $H_{WL}^{\text{imp}} = 20$ Oe impulse, are presented in Fig. 3.9.

---

5 To speed up computations we used higher value of $\alpha$ than experimentally measured 0.013 [99]
6 As mentioned in sec. 3.2.1 our samples have 1 nm TaN spacer. However, to get reasonable computation time we used 5 nm spacer.
Figure 3.8: A magnetization response (middle) of two 450 × 375 nm coupled layers to a 20 Oe amplitude WL impulse excitation (top), at $H_{EA} = 80$ Oe (a). The TJ resistance (bottom) in arbitrary units computed for each layer according to formula 3.12. Comparison of the SF critical field estimated from the micromagnetic simulations (open circles) and theoretical $H_{SF}$ (solid line) calculated from Eq. 3.2 as a function of the intrinsic anisotropy field (b).
3.3. RESULTS AND DISCUSSIONS

Figure 3.9: Results of the \( \mu \)-magnetic simulations for elliptical bi-layers with different lateral sizes. The magnetic parameters used in the simulations are given in the text. The top row is the TJ’s resistance \( R(t, H_{EA}) \) in response (for the 1\textsuperscript{st} magnetic layer) to a 200 ps 20 Oe impulse applied on WL (45\(^\circ\) to EA). The second (first layer) and third (second layer) rows are the FFT power spectral density color maps of \( R(t, H_{EA}) \).
The first row in figure 3.9 shows the $R(t, H_{EA})$ dependencies for one free layer. The second and third rows show PSD spectra for the first and the second layer, respectively. Due to the unequal field torques produced by the WL impulse on each layer (see Fig. 3.10), the resulting spectra for the two layers can be slightly different.

For $H_{EA} < H_{SF}$, i.e. when the moments are in the AP state, there are hardly any excitations in the spectrum. We have to keep in mind that in the AP$_{TJ}$ and P$_{TJ}$ states, the TJ resistance is rather insensitive to small-angle oscillations, as explained in sec. 3.2.2.2. This consideration is confirmed by our experiments, where the spectrum at $H_{EA} < H_{SF}$ is essentially flat.

The interlayer exchange, $J$, between the layers is set to zero in all simulations. The intrinsic anisotropy was set to 6 Oe in (c) and 5 Oe in Fig. 3.10(a,b,d). The value of $H_i$ does not affect the dynamics appreciably. However, a small change in the intrinsic anisotropy results in a significant change of the SF critical field. The comparison of the $H_{SF}$ estimated from the $\mu$MS and that calculated from the SDM (3.2) is given in Fig. 3.8(b). This is the first indication that the SD approximation does not model the spin dynamics in large junctions very well. Indeed, the spec-

![Diagram](image-url)
3.3. RESULTS AND DISCUSSIONS

The spectrum of the smallest particle in Fig. 3.9(a) exhibits almost perfect acoustical and optical modes having narrow linewidth. The linewidths become a few hundreds MHz for larger particles, suggesting that the oscillations of the magnetization contain a rather broad range of modes. This is an indication that the magnetic layers are behaving as nonsingle-domain particles. As we increase the size of the particles (b-d), the shape and the linewidth of the modes exhibit larger deviations from the single domain case, which can be summarized as follows. First, the micromagnetic spectra, \( f(H_{EA}) \), is more complex than the one predicted by the SDM (Fig. 3.6) at low and intermediate values of the magnetic field. According to the SDM, the optical branch goes down in frequency whereas the acoustical mode increases as the field is increased, forming a characteristic X-shaped spectrum – an optical-acoustical cross. In the spectra shown in Fig. 3.9(c,d) we can barely see this cross. However, the region past the cross point, where \( f_o \) decreases and \( f_a \) increases with field, is well recognized at high fields.

The second distinctive property of the micromagnetic spectrum is seen in Fig. 3.9(c) where the optical branch, \( f_o(H_{EA}) \), goes up with magnetic field after reaching a minimum at 170 Oe. This behavior is observed in our experiments for some TJs where the optical mode is pronounced. It was also observed in interlayer-exchange studies \[100\]. The SD model can also describe this upturn in frequency at high field if the expression for the optical frequency (3.7) is modified in the following way:

\[
    f_o \approx \frac{\gamma}{2\pi} \sqrt{4\pi M_s \left| H_{xsat} - \frac{H_{EA}}{H_{xsat}} \right|}. \tag{3.13}
\]

The third characteristic feature of the spectrograms shown in figures 3.9(c,d) is the presence of micromagnetic transitions at intermediate fields in the SS. These transitions are similar to the spin-flop abrupt transition, AP \( \Rightarrow \) SS, observed in \( R(t, H_{EA}) \) as a sudden appearance of oscillations at \( H_{SF} \). The transitions can be seen in Fig. 3.9(c,d-top row) as vertical lines in \( R(t, H_{EA}) \) plots. Some transitions are accompanied by a reversal of the oscillation phase, while in others the frequency of the oscillation shifts abruptly (Fig. 3.9(c,d) middle and bottom rows). To understand the physical mechanism responsible for these abrupt transitions we have performed a micromagnetic study of the spin distributions in the layers.

In the spectrogram for a 450 × 375 nm bi-layer shown in Fig. 3.9(c), we identify three micromagnetic transitions: the SF transition at \( H_{EA} = 48 \) Oe, and two additional transitions at 113 and 136 Oe. A detailed micromagnetic analysis of these transitions is illustrated in Fig. 3.11. Below the spectrogram (a) we plot angles of the two layers \( \theta_{1,2} \) between the EA and the averaged magnetization vectors in each layer (b). When magnetic field is increased from 0 to 44 Oe, the magnetic system remains in the AP configuration. The spin-flop transition at 48 Oe puts the moments into SS. Surprisingly, this SS is not symmetric with respect to the easy axis along which the magnetic field is applied. As the field is increased \( \theta_{1,2} \), shown as blue and red lines in Fig. 3.11(c), decreases continuously until 110 Oe, where an abrupt jump of the angle occurs. This is the first micromagnetic state transi-
Figure 3.11: Micromagnetic analysis of the spin-flop dynamics for a 450 × 375 nm TJ: power spectral density spectrum vs. $H_{EA}$ and $f$ (a); average angles of the two magnetic layers (red and blue) vs. $H_{EA}$ (b); 3D diagram of the averaged magnetization vectors illustrating the spin asymmetry of the SS (c); non-uniform spin distributions within the top and bottom ferromagnetic layers at critical $H_{EA}$ values corresponding to micromagnetic transition between the C-, and S-shaped states in the structure (d).

As the field is increased further, the angles again change continuously until 134 Oe where another transition takes place. This time, the moments jump into a symmetric scissored state with respect to the EA. When the field is increased further, the moments saturate symmetrically into the P state. The symmetry is shown in Fig. 3.11(b) by mirroring $-\theta_2$ (gray crosses), which coincides exactly with $\theta_1$ (red open circles) for $H_{EA} \geq 136$ Oe.

Such behavior cannot be explained within the SDM, where there is no ener-
getic reason for the two moments to be in a non-symmetric configuration with respect to EA, in the case of identical and symmetric coupled layers. We clarify this phenomena by studying the magnetization distribution in both layers, shown in Fig. 3.11(d), where the color map represents the in-plane $XY$ angles, $\theta_{1,2}$, between $M_{1,2}$ and EA. In the AF state ($H_{EA} \leq 44$), the magnetization is in a C-shaped state\(^7\). In the SS (SF at 48 Oe), the magnetization is distributed non-uniformly forming an S-shaped configuration noticeable already at 48 Oe. This S-state is well recognized at 108 Oe (3rd row of ellipses from the top in 3.11(d)). After the next transition (113 Oe) and the following increase of the field to 132 Oe, the magnetization is distributed non-symmetrically with respect to the EA and system jumps into C-shaped state again. The magnetization bending at the ends of both ellipses comes from the strong dipole interaction repelling the spins at the ends. Naturally, upon a further increase of $H_{EA}$ more and more spins gradually turn toward the EA direction (not shown in the figure). Here, after the last magnetization transition at 136 Oe, $f_{o,a}(H_{EA})$ dependencies are well defined as can be seen from the spectrogram in Fig. 3.11(a). For higher fields, when both moments are symmetric with respect to EA and magnetization forms S-state, we obtain a good agreement between the SDM and the $\mu$MS predictions.

In conclusion, we have studied the magnetization dynamics in spin-flop tunnel junctions numerically by using micromagnetic simulations. We obtain similar $f_{o,a}(H_{EA})$ dependencies from our $\mu$MSs and SDM in limited field ranges. $\mu$MS predict richer spectra, containing such features as micromagnetic state transitions at intermediate values of the magnetic field. At higher fields characteristic frequency modes are much better defined and are in good agreement with the SDM. The spectral linewidth is strongly affected (broadened) by non-uniform magnetization distributions in the layers. Our micromagnetic model shows that non-uniform spin distributions can be C- and S-shaped states. We also learned that the magnetization oscillations at fields below $H_{SF}$ are weak when detected through magnetoresistance, and so it can be difficult to detect them in typical experiments.

We have presented here the results of detailed micromagnetic analysis for 450×375 nm coupled layers only. Slightly different sizes and some variation in AR of the elliptical particles leads to similar results. We studied symmetrical particles in geometry and identical in their magnetic properties. The case of a thickness imbalance of the layers or a presence of an additional fringing magnetic field acting on one of the layers (e.g. field originating from the fixed layer) has not been studied, and remains a challenging and very interesting problem. We also do not present here the results of the $\mu$MS with non-zero exchange coupling between the layers, $J$. We note that the SF dynamics can be affected by the $J \neq 0$ in a non-trivial way. These effects, however, are outside of the scope of this work. Theoretical and experimental studies focused on dynamical properties of interlayer-exchange-coupled Ni$_{81}$Fe$_{19}$/Ru/Ni$_{81}$Fe$_{19}$ films have been recently published [100].

\(^7\)C-state is also known as "buckle" state [101].
The results of the μMS will help us to understand the experimental data presented in the next section.

### 3.3.2 Spin-flop dynamics: experimental results

The results presented in this section have been obtained using the measurement technique described in section 3.2.2.1. In order to excite the free layers in the SF TJ by magnetic field, we use short (200 ps halfwidth) impulses applied through the WL of the selected TJ cell. The amplitude of the impulse should not be too high or too low. We have tried different impulse amplitudes in the experiment. We learned that application of impulse fields larger than 30 Oe leads to appearance of additional features, outside the range for small signal oscillations of primary interest in this work. On the other hand, measurements with low impulse fields, < 10 Oe, give poor signal-to-noise (S/N) ratio. By experimenting we found the optimal impulse amplitude to be 20 Oe, which has been used in the data presented below. The measurements contained noise and possibly unidentified modes below 1.5 GHz. This noise did not depend on the impulse amplitude. Therefore, we conclude that it does not originate from the magnetic response of the TJ and we discuss the data in 1.5–6 GHz spectral range.

The measured voltage response $\Delta V_{TJ}(t, H_{EA})$ for a 490 $\times$ 400 TJ is shown in Fig. 3.12(b). The calculated PSD of these data is the spectrogram presented in Fig. 3.12(c). At low field, the optical branch $f_o$ is weak but is above the noise floor and has a constant frequency of about 4 GHz unto SF transition at $\sim$ 45 Oe. After the SF transition, the increase of the field leads to a decrease of $f_o$. At approximately 190 Oe the optical mode reaches its minimum frequency and, upon the following increase of the field, tends to go up in frequency again. This behavior is in good agreement with our micromagnetic model discussed in the previous section (see Fig. 3.9(c) and Eq. 3.13).

The characteristic of this SF junction is the presence of the micromagnetic transitions predicted by our micromagnetic model (sec. 3.3.1). Figure 3.12(a) clearly shows resistance jumps at approximately \{80; 145; 150\} Oe. The real time data (b) and spectrogram (c) also show micromagnetic transitions at the same critical fields. $f_o$, at these critical fields experience shifts and abrupt changes in intensity. This fact confirms our findings that at intermediate fields the spins are distributed non-uniformly across each free layer. This favors a non-symmetric configuration of the averaged magnetization vectors with respect to the EA direction, which can change abruptly at certain critical transition fields.

In Fig. 3.13 we compare experimental (a) and simulated (b, c) spectrograms for a 450 $\times$ 375 nm TJ. Unfortunately, the S/N ratio in this measurement is not very high. However, the $f_o(H_{EA})$ X-shaped dependence can be recognized and it agrees well with the micromagnetic model. Despite the high noise in the measurement, even the optical mode at $\approx$ 4 GHz and $H_{EA} < H_{SF}$ is clearly seen. We show two μMS spectra for comparison because the first one, having the $H_i = 2$ Oe, shows a weak optical mode in the AP state, while the second spectrum with
Figure 3.12: Measurement results for a 490 × 400 nm TJ. Static TJ resistance (a). The voltage response (b) is averaged 1024 times and is measured in mV. Spectrum (c) is the PSD of (b), in linear scale (a. u.). The bit is excited with a $H_{\text{WL}}^{\text{imp}} = 30$ Oe impulse, applied every 10 µs.

$H_i = 6$ Oe exhibits no optical mode in the AP state but a rather sharp SF transition at the same EA field as in the experiment (a). The main difference between the experiment and the micromagnetic results is that the measured acoustical branch seem to be split in two subbranches. This is not the case in the micromagnetic simulations where we assumed that both layers have the same magnetic properties and geometry. However, in the real TJs there might be present thickness imbalance leading to an unequal effective fields in free layers.

In conclusion, we studied the SF dynamics in TMRAM experimentally. The optical and acoustical modes, predicted by the single-domain model, have been detected in our experiments. We observed additional spectral features in excellent agreement with our micromagnetic simulations. The broadened linewidths and
Figure 3.13: Comparison of the experimental (a) and simulated spectra (b,c) for the 450x375 nm bit with following parameters: $H_{\text{imp}}^{\text{WL}} = 20$ Oe, 4096 averages in (a), intrinsic anisotropy is 2 and 6 Oe in (b,c) respectively. All three spectra are shown in linear scale. The AP part ($H_{\text{EA}} < 45$ Oe) of the spectrum in (a,b) is multiplied by factor of 2 for clarity.

Micromagnetic state transitions were measured and are explained using detailed micromagnetic simulations.
Chapter 4

Conclusions

In the first part of this work we studied STT effects in point contacts to single ferromagnetic films. We learned that, similar to F/N/F multilayers, single N/F interface can exhibit magneto-transport properties such as resistance excitations and even hysteretic magnetization reversal caused by nominally unpolarized currents. The experimental confirmation (paper I) of the predicted [61, 63] STT mechanism for single N/F interface helped us to understand the nature of the magnetic excitations (paper II) and the origin of the current-induced hysteretic switching (papers III, IV and [65]).

The investigations mentioned above involved static measurements on magnetic PCs. When looking at the progress of the theoretical and experimental research on STT in N/F multilayers or F/N/F spin-valve devices, we see that these studies were accompanied by investigations of STT-induced magnetization dynamics. It was an important and very tempting experiment (paper V) to probe spin dynamics at single N/F interfaces. By irradiating magnetic PCs with microwaves we observed a resonant stimulation of spin-wave modes. A direct detection of the magnetization oscillations at GHz frequencies in PCs to single F films remains a challenging experiment.

* * *

The second part of this thesis (paper VI) focused on the spin dynamics in spin-flop tunnel junctions used in toggle MRAM.

The single-domain model of two coupled nanomagnets [95] predicts two characteristic resonant modes. Our real-time measurements on spin-flop TJs, excited by 100 ps rise time impulse fields, confirmed this prediction. In addition, we found that the spectrum of the observed GHz oscillations contains magnetic transitions which we explain with the help of micromagnetic simulations.

Our results should be useful for understanding the limitations on the operating speed of toggle magnetic tunnel junctions and designing new types of related magnetic storage elements.
Bibliography


[95] D. C. Worledge. UNPUBLISHED. Private communication with author. Numerical and analytical results of the spin-flop dynamics within the single-domain approximation will be published elsewhere.


Chapter 5
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«Try not to become a man of success,
but rather try to become a man of value.»
Albert Einstein

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Chapter 6

Appended papers

6.1 List of papers

Author contribution: fabrication of samples.

Author contribution: fabricated the samples, conducted most of the measurements, wrote the manuscript.

Author contribution: fabricated the samples, conducted some of the measurements, took part in preparing the manuscript.

Author contribution: fabricated the samples, conducted the measurements, performed the micromagnetic simulations, wrote the manuscript.

V O. P. Balkashin, V. V. Fisun, I. K. Yanson, L. Yu. Triputin, A. Konovalenko, V. Korenivski, Spin dynamics in point contacts to single ferromagnetic films, Submitted to Phys. Rev. B.
Author contribution: fabricated the samples, some of the measurements conducted.

**Author contribution:** performed the measurements measurements, numerical data analysis and micromagnetic simulations, wrote the manuscript.

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1The single-domain theory was performed by D. C. Worledge.