Chapter 2

Magnetization switching dynamics
2.1 Introduction

Today, magnetism is the base for most of the data storage devices. In magnetic memory devices, logical bits (ones and zeros) are stored in a unit cell of magnetic material by setting the orientation of the magnetization vector inside the cell in one of two possible directions either up or down, coding the logical values 0 and 1. To write data, an external magnetic field is applied to reverse the magnetization, flipping the bit between 0 and 1. It is generally accepted that the fastest way to record a bit is to reverse the magnetization via magnetic field-induced precessional motion. The study of magnetization processes in magnetic materials has been the focus of considerable research for its application to magnetic recording technology in the last fifty years. In fact, nowadays the design of widespread magnetic storage devices, such as the hard-disks within computers, requires the knowledge of the microscopic phenomena occurring within the magnetic media. In this respect, it is known that some materials, referred to as ferromagnetic materials that has spontaneous magnetization at room temperature which is the result of spontaneous alignment of the elementary magnetic moments that constitute the medium. Roughly speaking, from a phenomenological point of view, one has a medium whose magnetization state can be changed by means of appropriate external magnetic fields. The magnetic recording technology exploits the magnetization of ferromagnetic media to store information.

In 1935 Landau and Lifshitz [175] proposed an equation of motion for magnetization in a homogeneously magnetized body. This equation was modified by Gilbert [216] to overcome the unphysical solution for large damping parameters [217]. The so-called Landau-Lifshitz-Gilbert (LLG) equation consists of two terms. The first term represents the Larmor precession of magnetization about the direction of the internal field and the second term is a phenomenological damping term that describes the energy dissipation of the system. The Gilbert damping constant controls the dissipation rate associated with either the small amplitude motion probed in ferromagnetic resonance or Brillouin light scattering studies of long wavelength spin excitations in ferromagnetic media and also
that of large amplitude spin motion associated with magnetization reversal. As
the LLG equation is a nonlinear partial differential equation, analytical solu-
tions can be found only in special cases. The switching process of a magnetic
element sensitively depends on intrinsic properties, on the physical or chemical
microstructure, and on the characteristics of the external field pulse. The first
study of magnetization switching using the LLG equation was performed by
Kikuchi [218]. An analytical solution was found for the magnetization switch-
ing of an isolated, isotropic, single domain sphere. In 1958, anisotropy was
included in this calculation [219]. However, detailed solutions of the LLG equa-
tion require extensive numerical integration [220-222].

The numerical solution of the Gilbert equation of motion provides the theo-
retical background for the switching process in ferromagnetic structures. The
switching time considerably depends on the Gilbert damping constant $\alpha$. Kikuchi
[218] calculated the minimum reversal time of a single domain sphere and
a single domain thin film with in-plane anisotropy. The critical value of the
damping constant that minimizes switching time is $\alpha = 1$ for a sphere and
$\alpha = 0.01$ for a thin film. He and Doyle [221] solved the Landau-Lifshitz equation
numerically to investigate switching with very short field pulses. They conclude
that switching time of the order of about 100 ps are possible if the external field
is applied at 90° with respect to the anisotropic axis. Bauer and co-workers
[223] investigated the switching properties of magnetic thin-film element sub-
ject to ultrashort magnetic field pulses numerically. The thin film is described
by Stoner-like magnetic blocks. Field pulses of 4 ps duration cause magne-
tization reversal in both perpendicular and in-plane magnetized films. How-
ever, the time needed for magnetization to reach equilibrium is about 550 ps.
Mallinson [224] derived switching time as a function of field strength for fields
parallel to the anisotropy direction. Switching time decreases with increasing
external field. Koch and co-workers [225] investigated the switching dynam-
ics of microsized magnetic thin films experimentally and numerically. They
observed switching time which is below 500 ps. Albuquerque and co-workers
[226] presented a finite difference method to solve the Gilbert equation effec-
tively for thin film structures used in current-tunnel junction MRAM devices. A sequence of tailored field pulse causes quasi-coherent switching in the sub-nanosecond regime.

2.2 Magnetization Dynamics

In this section, some insights in the origin of the precessional spin motion are offered and various relaxation dynamics are discussed. It is known from quantum mechanics that there is a proportionality relationship between the magnetic spin momentum $\vec{\mu}$ and angular momentum $\vec{L}$ of electrons. This relationship can be expressed as

$$\vec{\mu} = -\gamma \vec{L},$$

(2.1)

where $\gamma = 2.21 \times 10^5 mA^{-1}s^{-1}$ is the absolute value of the gyromagnetic ratio and $\gamma = g q \mu_0 / 2 m_e$, here $g \sim 2$ is Lande’s splitting factor, and $q$ and $m_e$ represent the charge and mass of the electron respectively. When a magnetic moment $\vec{\mu}$ is placed in a magnetic field at an angle, it will experience a torque $\vec{\tau} = \vec{\mu} \times \vec{H}$, because the magnetic moment has an angular momentum, the torque will cause the angular momentum to precess about the magnetic field, according to Newton’s classical equation of motion, $\frac{d\vec{L}}{dt} = \vec{\tau}$ is given by

$$\frac{d\vec{L}}{dt} = \vec{\tau} = \vec{\mu} \times \vec{H}.$$

(2.2)

The relation between the magnetic moment and the angular momentum is written in terms of the gyromagnetic ratio $\gamma$ according to Eq. (2.1) and the equation of motion is given by

$$\frac{d\vec{\mu}}{dt} = -\gamma \vec{\omega} = -\gamma [\vec{\mu} \times \vec{H}].$$

(2.3)

This leads to the Larmor precession of magnetic moment $\vec{\mu}$ about the magnetic field $\vec{H}$ at the angular frequency $\omega = -\gamma \vec{H}$. In a ferromagnetic material, the atomic spins are coupled together by very strong exchange interaction of microscopic origin. As a consequence, if the system is driven by an externally
applied magnetic field or by some other perturbation that drives the magnetization away from its equilibrium orientation and if the spins are examined within a very small volume \(dV_\tau\), the spins remain locked tightly parallel to each other by virtue of exchange. A consequence is that the system may be described completely by its magnetization per unit volume \(\vec{S}(r, t)\), a vector of fixed length as it precesses. Therefore, Eq. (2.3) can be written for each spin magnetic moment within the elementary volume \(dV_\tau\)

\[
\frac{d\vec{\mu}_j}{dt} = -\gamma[\vec{\mu}_j \times \vec{H}],
\]

(2.4)

where the magnetic field \(\vec{H}\) is intended to be spatially uniform. Now, by taking the volume average on both sides, one can obtain

\[
\frac{1}{dV_\tau} \frac{d}{dt} \sum_j \vec{\mu}_j = -\gamma \frac{\sum_j \vec{\mu}_j}{dV_\tau} \times \vec{H}.
\]

(2.5)

By using the definition of magnetization vector field \(\vec{S} = \frac{\sum_j \vec{\mu}_j}{dV_\tau}\) in Eq. (2.5), we obtain the following equation for precessional motion

\[
\frac{\partial \vec{S}}{\partial t} = -\gamma[\vec{S} \times \vec{H}].
\]

(2.6)

Figure 2.1: Precession of the magnetic moment about the magnetic field \(\vec{H}\) direction.

Fig. (2.1) shows the precession of the magnetic moment of an electron about the field direction. The motion of \(\vec{S}\) is in agreement with the conservation of angular momentum. Because of Newtons law action=reaction, the torque acting on \(\vec{S}\) acts with the opposite sign on the source of the magnetic field \(\vec{H}\) and, in principle, sets it in opposite rotation so that the total angular momentum is
conserved. Eq. (2.6) represents the precessional motion which can be theoretically justified either on quantum or classical grounds. By considering a static particle with spin in an uniform magnetic field, the time evolution for the mean value of the spin operator can be derived using either Schrödinger equation [227] or equivalently, Von Neumann equation [216]. On the classical ground, Eq. (2.6) can be justified by assuming that the magnetic moment arises from a “circular” motion of an electron and by using Newton’s law for the angular momentum, as well as the relation between the magnetic moment and angular momentum [228,229]. The idea of this classical explanation was first given by Sir Joseph Larmor and that is why the precessional motion of the magnetic moment about the magnetic field is often called as Larmor precession. Since in 1 T field the spin precesses by 1 rad = 360°/2π = 57.3° in 5.7 ps, it is convenient to remember that it takes about 10 ps for a 90° spin precession, which is required for initiating a switch of the magnetization into the opposite direction. The magnetization dynamics in ferromagnets caused by typical external field occurs on the tens of picosecond time scale. With the advent of spectroscopic techniques to measure the magnetization in a time as short as \( \sim 10^{-12} \) s, the precession can now be directly imaged in the time domain.

### 2.3 Landau-Lifshitz equation

The first dynamical model for the precessional motion of the magnetization was proposed by Landau and Lifshitz. Basically, this model is constituted by a continuum precession Eq. (2.6), in which the presence of quantum mechanical effects and anisotropy is phenomenologically taken into account by means of the effective field \( \vec{H}_{\text{eff}} \). Then, the Landau-Lifshitz (LL) equation is given by

\[
\frac{\partial \vec{S}}{\partial t} = -\gamma [\vec{S} \times \vec{H}_{\text{eff}}],
\]  

(2.7)

where the effective field \( \vec{H}_{\text{eff}} \) comprises both external and internal field acting on the magnetization: \( \vec{H}_{\text{eff}} = \vec{H}_{\text{stat}} + \vec{H}_{\text{pulse}} + \vec{H}_{\text{shape}} + \vec{H}_{\text{ani}} \) with \( \vec{H}_{\text{stat}} \) the applied static field, \( \vec{H}_{\text{pulse}} \) the magnetic field pulse, \( \vec{H}_{\text{shape}} \) the shape anisotropy field, and \( \vec{H}_{\text{ani}} \) the sum of all other anisotropy fields, which in general include mag-
netocrystalline, magnetoelastic and interfacial anisotropy contributions. It is found that if the rate of change of magnetization $\frac{d\mathbf{S}}{dt}$ vanishes in Eq. (2.7), then it expresses the equilibrium condition given by the Brown’s equation $\gamma [\mathbf{S} \times \mathbf{H}_{eff}] = 0$. The Brown’s equation allows us to find the equilibrium configuration of the magnetization within the body and it states that the torque exerted on magnetization by the effective field must vanish at the equilibrium. It may be noted that LL equation is a conservative equation, but dissipative processes which takes place within the dynamical system. Now, we consider the relaxation motion of the magnetic moment, which can be quantitatively described by adding a dissipative (damping) correction term to Eq. (2.7). The dissipated energy is actually transformed by various mechanisms into the thermal energy of a system. Although these mechanisms are partially known [230,231], they are too complex to be taken into account in an explicit derivation of the damping correction term at a macroscopic level. In order to describe the experimental results, various phenomenological expressions are employed. Most notable one were given by Landau and Lifshitz for the description of energy losses in the magnetic domain wall motion in ferromagnetic materials [175], and by Bloch for the description of nuclear magnetic relaxation [232]. Since these expressions have been successfully applied to various physical phenomena involving dissipation of the magnetic energy. The Landau-Lifshitz expression is mostly used for the description of various dissipative processes in which the norm of the magnetic moment is conserved, while the Bloch expression is appropriate to complementary cases. Then, the Landau-Lifshitz equation becomes

$$\frac{d\mathbf{S}}{dt} = -\gamma [\mathbf{S} \times \mathbf{H}_{eff}] - \frac{\lambda}{M_s} [\mathbf{S} \times \mathbf{S} \times \mathbf{H}_{eff}],$$

(2.8)

where $\lambda > 0$ is a phenomenological characteristic constant of the material and this additional term pushes the magnetization in the direction of $\mathbf{H}_{eff}$. The microscopic origin of the magnetic damping constant is a manifold. Possible mechanisms are direct coupling to the lattice via spin-orbit interaction or indirect coupling via spin waves, e.g., two or higher order magnon scattering processes [233] or impurity relaxation mechanisms [230,234]. In conducting
materials, free electrons add an additional contribution via eddy currents to the magnetic damping.

2.4 Landau-Lifshitz-Gilbert equation

When the magnetic moment is placed in an external static magnetic field, the magnetic moment eventually moves into the field direction. This fact underlies the earliest magnetic device, the compass. We have seen that the precessional torque \( \vec{\tau} = \vec{S} \times \vec{H}_{eff} \) cannot accomplish this situation as it is perpendicular to \( \vec{H}_{eff} \). From this observation, it may be noted that an additional torque \( \vec{\tau}_D \) must be introduced which is perpendicular to both the precessional torque and the magnetization to describe a dissipation phenomenon. We shall call \( \vec{\tau}_D \) the damping torque as proposed by Gilbert [222],

\[
\vec{\tau}_D = C \left[ \vec{S} \times \frac{\partial \vec{S}}{\partial t} \right].
\] (2.9)

The constant of proportionality \( C \) is purely phenomenological, analogous to the friction coefficient for a linear motion. According to Fig. (2.2), \( \vec{\tau}_D \) causes a rotation toward the +z axis for positive \( C \). Therefore, the precessional Eq. (2.7), modified according to Gilbert's work, is generally referred to as Landau-
Lifshitz-Gilbert (LLG) equation

\[
\frac{\partial \vec{S}}{\partial t} = -\gamma [\vec{S} \times \vec{H}_{\text{eff}}] + \frac{\alpha}{M_s} \left[ \vec{S} \times \frac{\partial \vec{S}}{\partial t} \right].
\] (2.10)

The phenomenological constant \(\alpha\), so called damping parameter, stands for unspecified dissipation phenomena. The first term describes the precession of \(\vec{S}\) at a fixed angle \(\theta\) about the magnetic field direction \(\vec{H}_{\text{eff}}\), as illustrated in Figs. (2.1) and (2.2) due to the torque \(\vec{\tau} = \vec{S} \times \vec{H}_{\text{eff}}\) into the direction \(\frac{\partial \vec{S}}{\partial t} = \gamma \vec{\tau}\). The second term describes the change of \(\vec{S}\) due to the damping torque \(\vec{\tau}_D\), causing \(\vec{S}\) to turn toward the direction of \(\vec{H}_{\text{eff}}\) for a positive value of \(\alpha\), as shown in Fig. (2.2). Solving the LL or the LLG equation shows how \(\vec{S}\) spirals into the magnetic field direction as depicted in Fig. (2.3(i)). There is substantial

![Figure 2.3: Precessional motion of the magnetic moment \(\vec{S}\) according to (2.10) (i) for positive damping parameter and (ii) negative damping parameter. (i) \(\vec{S}\) spirals into the direction of the field \(\vec{H}_{\text{eff}}\) at which point both the precessional torque and the damping torque vanish. (ii) \(\vec{S}\) spirals away from the direction of the field \(\vec{H}_{\text{eff}}\) into a final position opposite to \(\vec{H}_{\text{eff}}\) thereby gaining energy.](image)

difference between LL and LLG equations although they are very similar from mathematical point of view. For instance, LL Eq. (2.8) can be obtained easily from LLG equation. In fact, by vector multiplying both sides of Eq. (2.10) by \(\vec{S}\), one obtains

\[
\vec{S} \times \frac{\partial \vec{S}}{\partial t} = -\gamma \vec{S} \times (\vec{S} \times \vec{H}_{\text{eff}}) + \vec{S} \times \left[ \frac{\alpha}{M_s} \vec{S} \times \frac{\partial \vec{S}}{\partial t} \right].
\] (2.11)
By using the vector identity in Eq. (2.11), we get
\[ \vec{S} \times \frac{\partial \vec{S}}{\partial t} = -\gamma \vec{S} \times (\vec{S} \times \vec{H}_{\text{eff}}) - \alpha M_s \frac{\partial \vec{S}}{\partial t}. \] (2.12)

By substituting Eq. (2.12) in the right hand side of LLG Eq. (2.10), one has
\[ \vec{S} \times \frac{\partial \vec{S}}{\partial t} = -\gamma \vec{S} \times \vec{H}_{\text{eff}} - \frac{\gamma \alpha}{M_s} \left[ \vec{S} \times (\vec{S} \times \vec{H}_{\text{eff}}) \right] - \alpha^2 \frac{\partial \vec{S}}{\partial t}, \] (2.13)

which can appropriately be written as,
\[ \vec{S} \times \frac{\partial \vec{S}}{\partial t} = -\frac{\gamma}{1 + \alpha^2} \vec{S} \times \vec{H}_{\text{eff}} - \frac{\gamma \alpha}{(1 + \alpha^2) M_s} \left[ \vec{S} \times (\vec{S} \times \vec{H}_{\text{eff}}) \right], \] (2.14)

which is commonly referred to as Landau-Lifshitz equation in the Gilbert form.

One can immediately notice that Eq. (2.14) and Eq. (2.8) are mathematically the same, provided that one assumes,
\[ \gamma_L = \frac{\gamma}{1 + \alpha^2}, \quad \lambda = \frac{\gamma \alpha}{1 + \alpha^2}. \] (2.15)

Moreover, Podio-Guidugli [235] has pointed out that both LL and LLG equations belong to the same family of damped gyromagnetic precession equations. Nevertheless some considerations about the quantity \( \gamma \), which indeed is the ratio between physical characteristics of the electrons like mass and charge, are sufficient to say that Eqs. (2.8) and (2.10) express different physics and are identical only in the limit of vanishing damping. Further, first Kikuchi [218] and then Mallinson [217] have pointed out that in the limit of infinite damping (\( \lambda \to \infty \) in Eq. (2.8), \( \alpha \to \infty \) in Eq. (2.10)), the LL equation and the LLG equation give respectively,
\[ \frac{\partial \vec{S}}{\partial t} \to \infty, \quad \frac{\partial \vec{S}}{\partial t} \to 0. \] (2.16)

Since the second result is in agreement with the fact that a very large damping should produce a very slow motion while the first is not, one may conclude that the LLG equation is more appropriate to describe magnetization dynamics. The damping may be thought of as being generated by an effective magnetic field \( \vec{H}_{\text{eff}} \propto \frac{\partial \vec{S}}{\partial t} \). Induction current indeed create a magnetic field varying with the frequency of the precessing magnetization \( \vec{S} \) but in nanoscopic materials and
thin films they cannot contribute to the damping torque $\vec{r}_D$ [236]. Therefore, different mechanisms than that of electromagnetic induction must be active in thin films to explain the observed damping. As shown in Fig. (2.3(i)), during the damping process the magnetization spirals into the field direction and this process corresponds to a change in angular momentum. Angular momentum conservation then demands that during the damping process angular momentum must be transferred to another reservoir. Initially angular momentum may be transferred within the spin system itself by excitation of spin waves, but ultimately the angular momentum is transferred to the lattice. Therefore the LL and the LLG equations describe the temporal evolution of a magnetization $\vec{S} = \vec{u}/V$, assuming that the magnitude $|\vec{S}|$ remain constant in time. In 1996 John Slonczewski [237] and Luc Berger [238] independently proposed that the damping torque may have a negative sign as well, corresponding to a negative sign of $\alpha$. Under this condition, $\vec{S}$ moves into a final position anti-parallel to $\vec{H}_{eff}$ as illustrated in Fig. (2.3(ii)). In magnetism, the negative damping phenomenon is highly appreciated since it allows one to switch the magnetization. Clearly, energy has to be supplied to achieve the motion of $\vec{S}$ in which the angle $\theta$ between $\vec{S}$ and $\vec{H}_{eff}$ is enlarged. This energy is thought to be provided by injecting spin-polarized electrons from an adjacent ferromagnet, magnetized in the opposite direction compared to the magnetic material under consideration. The injected electrons are the minority spins and therefore has higher energy than the average electrons. If their spin polarization is conserved, they will add to the magnetization component anti-parallel to $\vec{H}_{eff}$ thus effectively enlarging $\theta$. Pulsed magneto-optics in the visible and X-ray energy range offers the possibility to directly observe the motion of $\vec{S}$.

2.5 Switching of magnetization

Following the huge development of electronic devices such as the personal computer, this era has been strongly driven by the growing demand to increase the density and the speed of writing and retrieving data in memory devices. Today, the demand for information storage is enormous and expected to increase
even further as new technology arrives. We will hereafter focus our attention on the magnetization reversal processes. Magnetic reversal or switching is the process by which the magnetization of a specimen is changed from one stable direction into another. In practice, it involves a rotation of the magnetization by 180°, from one orientation along the easy axis to the opposite orientation, and this process is therefore referred to as magnetization reversal. Technologically, this is one of the most important processes in magnetism that is linked to the magnetic data storage process. It is important to realize that in today's technological applications the magnetization is typically not read immediately after its switching is initiated or completed. Therefore in technology, the relevant time for the writing process is only the time that it takes to put the system into a state from which it will reliably move into the new desired magnetization direction. As it is known today, there are only few possible ways to reverse the magnetization of a metallic magnet: reversal through an applied magnetic field, reversal by optical pulse, reversal by spin injection, the magnetization reversal by means of stress-induced anisotropy and by thermal activation. At the outset, we start with the discussion of magnetization switching of magnetic materials through an applied magnetic field.

2.5.1 Magnetization switching in the presence of an applied magnetic field

Based on experimental advances, magnetization reversal has undergone considerable development in recent years [239-241]. More interestingly we note that the reverse of magnetization \( \vec{S} \) is determined as a function of the angle at which the external magnetic field \( \vec{H} \) is applied to the particle. The reversal mechanism is difficult to understand in detail, because \( \vec{S} \) can assume complex curling and buckling modes depending on the details of the shape and magnetic properties of the particle. In this section, we report some recent results regarding two different ways to achieve magnetization switching; the so called damping switching and precessional switching.
Damping Switching

Damping switching is a phenomenon where the applied field is reversed to switch the magnetization back to its initial orientation. When the field is switched in a time equal to half a precession period, the magnetization switches in the opposite direction. This process is known as damping switching. The figure illustrates this process, showing the initial orientation of the magnetization and the reversed orientation after damping switching.

2.5 Switching of magnetization

2.5.1 A

Damping Switching
2.5.1.B Precessional switching

The fastest and most economical switching is the precessional switching. The precessional switching is a new strategy to realize magnetization reversal which has been of considerable interest to the researchers in the recent years [242,243]. In the precessional switching, the magnetic field is applied in the direction perpendicular to the initial magnetization in order to use the associated torque to control magnetization precessional motion as depicted in Fig. (2.4b). In fact, this torque pushes the magnetization out-of-plane, creating a strong demagnetizing field in the direction perpendicular to the film plane. Then the magnetization starts to precess around the demagnetizing field. After the field pulse vanishes, the magnetization vector precesses like spinning-top around the demagnetizing-field direction before settling into the switched state. The reversal is obtained after half precessional oscillation. This kind of switching is much faster and it requires lower applied field with respect to the traditional switching. Moreover, in this scheme the magnetization may be switched back without changing the polarity of the magnetic field simply by applying the switching pulse again. Therefore, the only ingredients required to control the precessional switching are the strength and the duration of the pulses. The precessional switching process consists of two stages; in the first stage the magnetization precesses under the influence of an applied external field until its orientation is almost reversed, in the second stage the external field is switched off and the magnetization undergoes relaxation oscillation towards the nearby equilibrium point. In the first part of the process, the magnetization dynamics is typically so fast that dissipative effects can be neglected. On the other hand, dissipation has to be taken into account during the relaxation process. But it was realized that if an instantaneous magnetic-field pulse was used instead, reversal could happen more quickly by precession.

2.5.2 Switching by spin injection

It has been recently shown, both theoretically and experimentally, that a spin polarized current when passing through a small magnetic conductor can
current is reversed (now entering through the free layer) a spin accumulation in the middle of the magnetic layer. If the direction of the injected random spin polarization is reversed, the second layer, which is magnetized in the opposite direction, creates a negative magnetic field that opposes the injected polarization. The net result is a transverse magnetic field, which is perpendicular to the plane of the magnetic layers. This field induces a spin accumulation that is transverse to the injected polarization. The transverse field is generated by the magnetization of the second magnetic layer, which is magnetized in the opposite direction of the first layer. The net result is a spin accumulation that is transverse to the injected polarization. This effect is used in magnetic field sensors, where the transverse field causes a change in the sensor output.
Several studies have focused on magnetization pulse induced switching of the macro-
the LLC equation with an additional term [246]. From a different point of view,
in the behavior of the average spin magnetization vector can be described by
the quantum theory of the phenomenon is fairly well understood. Interest-
scope phenomena have been continuously confirmed [246, 247]. Although, the micro-
many contributions have been conducted on this geometry and the
information processing application.

Magnetization reversal technique is expected to lead to future data-storage and
magnetization beak to the original state as illustrated in the third Fig. (2.5). This new
spins are oriented opposite to the transmitted spins and can switch the magnet-
builds up in front of the polizer by reflection of spins from the polizer. These

The sensor back to the original state.

Spin have the opposite direction from the transmitted spins and can switch
front of the polizer through reflection of spins on the polizer. The reflected
When the current direction is reversed, a spin accumulation builds up in
milled current can switch the sensor layer by spin torque as shown in the mid-
though it located a second ferromagnetic sensor layer, the spin polarized cores-

Figure 2.5: Schematic illustration of a spin injection structure and the as-
experimental studies have also focused on spin current induced switching in the presence of a magnetic field, switching behavior for different choices of the angle of the applied field, variation in the switching time, etc., [249,250]. A numerical study on the switching phenomenon induced by a spin current in the presence of a magnetic field pulse has also been investigated very recently in [251]. As an extension to two dimensional spin configurations, the switching behavior on a vortex has been studied in [252].

### 2.5.3 Reversal by thermal activation

In this section, we shall discuss the reversal process under the influence of thermal energy. Magnetization reversal is simply the reversal of the direction of the magnetic moment of a magnetized sample caused by applying a magnetic field in the direction opposite to the moment. Most of the existing studies on the magnetization switching phenomenon have examined the situation wherein a pulsed magnetic field greatly exceeding the value of the zero-temperature switching field $H_C$ ($T = 0$) is applied, whereupon the switching of magnetization appears to begin immediately and finishes in a few nanoseconds or less. An equally important regime for magnetization reversal is one wherein the applied field is less than $H_C$. In the classical picture at $T = 0$, the magnetization does not reverse the field $H < H_C$, but remains stuck in a metastable local energy minimum and the reversal requires overcoming a barrier $\Delta E$ between adjacent minima by applying an external magnetic field. At $T > 0$, the thermal fluctuations help to overcome $\Delta E$ and thus magnetization reversal becomes temperature assisted. As it is presented in Fig. (2.6) the energy of uniaxial anisotropic magnetic particle has two minima and the magnetization orientation lies initially into one energy minimum. Thermally-assisted magnetization reversal has first been considered by Néel and Brown [253,254] and so-called Néel-Brown model. According to the fundamental theory of fine ferromagnetic particle theory proved by Brown, there is a critical diameter, below which the nanoparticle is uniformly magnetized and the magnetization reversal occurs by uniform mode. For somewhat larger ones, nonlinear reversal
modes (such as curling mode and buckling mode) are expected but the particle can still be considered as a monodomain; for even larger samples the magnetization has a multidomain structure and magnetization reversal may occur via domain wall propagation process [50,255]. This process has been studied on length scales ranging from molecular [256] to macroscopic, where it is connected with magnetization creep [257]. For the past few decades, the number of experiments has been performed on the nano particles to verify the Néel-Brown model for thermal activation of single-domain ferromagnets over an energy barrier [239,240]. According to the Stoner-Wohlfarth model [258], the energy barrier of a single-domain ferromagnetic nanoparticle in a field applied along the anisotropy direction can be written as $\Delta E = \Delta E_0 (1 - H/H_{SW}^0)^2$, where $\Delta E_0$ is the anisotropy energy barrier in zero field and $H_{SW}^0$, the field for which the energy barrier vanishes, is the switching field at 0 K. Street and Wooley have postulated an exponential decay of the magnetization from the metastable states with a relaxation time (a characteristic wait time) obeying the Arrhenius-Neel law $1/\tau = f_0 \exp(-\Delta E/KT)$ [259,260], where the pre-exponential factor $f_0 = 1/\tau_0$ is often taken as the Larmor precession frequency of the order of

![Diagram](image.png)

Figure 2.6: Thermally activated switching. Variation of energy as a function of angle $\theta$ between the direction of magnetization and easy anisotropy axis.
$10^9 - 10^{10}$ Hz and depends on variables like anisotropy constant, magnetization and damping [261,262]. Recent experiments indicate values of $\tau_0 = 4 \times 10^{-9}$ s for small ferromagnetic particles [48,26] and $\tau_0 = 1 \times 10^{-9}$ s for particular magnetic recording media [263]. The aforementioned models are valid for strictly uniform magnetization reversal. Usually, the reversal field can be lowered when the magnetization is not uniform during the reversal. Frei et al., [264] calculated that for a prolate ellipsoid the critical size shows little dependence on magnetocrystalline anisotropy and the exact shape is approximately equal to $A^{1/2}/M_S$, where $A$ is the exchange constant. For slightly larger particles, the magnetization will stop being uniform, and can arrange according to different spatial distributions.

2.5.4 Reversal through stress induced anisotropy

Possibilities of reversing magnetization without applying external magnetic field have been attracting researchers attention. The magnetic behavior of rigid ferromagnetic materials has been studied extensively during the last four decades on the basis of micromagnetic theory [73]. The coupled magnetoelastic phenomena are essential for the construction of efficient storage devices and sensor devices. The theoretical framework for the description of such coupled magneto-mechanical phenomena was proposed by Brown [265] and an extensive literature is cited in Refs. [266,267]. The foundation of such an approach is due to the pioneering works of Tiersten [268,269], Brown [50,265] and Maugin and Eringen [270-272]. Among the various phenomena that are present, magnetostriction is the most promising for applications and still the most difficult to be described theoretically. Magnetostriction is a property of ferromagnetic materials that causes them to change their shape or dimensions when subjected to a magnetic field. The effect was first identified in 1842 by James Joule when observing a sample of nickel. The term magnetostriction, in the classical sense, denotes the deformation of a ferromagnetic crystal when it is cooled under the critical Curie temperature, or when a previous saturated ferromagnetic crystal is applied an additional external field capable of increasing the spon-
taneous magnetization beyond its saturation value (forced magnetostriction). Whether the ferromagnetic crystal is lengthened or shortened in the direction of the magnetization, the material is characterized as having positive or negative magnetostriction, respectively. The inverse effect is also possible. When a ferromagnetic specimen is under mechanical loading (tension or compression) a change is noticed in its magnetic structure called inverse magnetostrictive effect. The capability of an applied mechanical loading to produce net resultant magnetization in a previously demagnetized ferromagnetic specimen has been examined experimentally and theoretically by Misra [273]. A wide class of phenomena that relates the elastic properties of ferromagnetic materials with the magnetization is generally referred to as magnetostriction (Joule effect, Matteucci effect, Wiedemann effect, etc.,) [274]. Among various approaches, the

![Diagram](image)

**Figure 2.7:** Stress-driven magnetization process. The direction of magnetization and the state of strain in the film plane for succeeding steps in the process.

magnetization switching by means of stress-induced anisotropy seems promising for electronic applications. A magnetostrictive film with 4-fold magnetocrystalline anisotropy is proposed as a candidate with which full reversal can be performed by a cycle of uniaxial stress/strain introduced along a fixed direction in the film plane [275] as illustrated in Fig. (2.7). In this method, initially
the magnetization vector resting in one of the easy axes pointing almost parallel to x-direction. A tension is given in the y-direction as the input to the system, which creates an easy axis of stress-induced anisotropy in y-direction. When the stress-induced anisotropy becomes strong enough, the magnetization vector \( \vec{S} \) rotates to align parallel to the easy axis. At this occasion of rotation, it should be noted that \( \vec{S} \) does not rotate towards -y but towards +y direction (Fig. 2.7b) because \( \vec{S} \) was initially tilted towards +y. Then the \( \vec{S} \) vector relaxes into the nearest easy axis of the 4-fold anisotropy as the stress induced anisotropy is removed (Fig. 2.7c). The second input to the system is compression. While the direction of this strain is again parallel to the y-axis, the easy axis of stress-induced anisotropy arises in x-direction due to the different sign of the strain. This stress-induced anisotropy rotates the magnetization vector towards the negative x-direction as illustrated in the first panel of Fig. (2.7d). As it can be seen from the following Fig. (2.7e), the cycle of compression and expansion applied along the y-axis resulted in successive 90° counter clockwise rotations of the magnetization vector. This type of driving has advantages that one does not need to drag the magnetization vector continuously over the whole range of angle between 0° and 90° or 180°. In this method, low power consumption is expected in the device because it would not require high current to produce magnetic field for switching.

### 2.5.5 All optical switching

Optical pulses could serve as an alternative stimulus to trigger magnetization reversal. In this section, all optical switching of metallic magnets is demonstrated as one of the possible ways of reversing the magnetization. Generally called all optical switching, this technique refers to a method of reversing magnetization in a ferromagnet simply by circularly polarized light where the magnetization direction is controlled by the light helicity. In fact, this switching process could be seen as similar to magnetization reversal by spin injection. The only difference is that the angular momentum is supplied by the circularly polarized photons instead of the polarized electrons. This hypothetical method
first proposed by Hübner and collaborators [276] is based on the application of a well-shaped ultrashort laser pulse of suitable frequency, polarization and duration and it potentially works on the femtosecond time scale. It involves the excitation of the system from the ground state to a well-defined final state where the magnetization is manipulated so that it decays into a new ground state with opposite magnetization direction as illustrated in Fig. (2.8). In this experiment, the laser pulse has two effects on the magnetization. First, it rapidly pumps energy into the film, locally heating the material and demagnetizing it [277]. The

![Circularly Polarized Laser](image)

Figure 2.8: The faster switching mode is driven by a short and intense circularly polarized laser pulse. Shortly after the pulse, the magnetization collapses due to heating. During the cooling, $\tilde{S}$ restores along the effective magnetic field $H_{IFE}$ that is generated by the inverse Faraday effect.

energy of the laser pulse is primarily absorbed by the electrons, which reach a temperature of about 1200 K within the first few hundred femtosecond after the pulse. Changes in the electronic temperature affect the magnetic properties in sub picosecond time scale. Most importantly, the magnitude of the magnetization $\tilde{S}$ decreases as the temperature of the electronic system approaches
the Curie temperature $T_C$. They show that the magnetization can be temporarily destroyed down to a value of zero about 500 $fs$ after applying a sufficiently strong laser pulse. They used the well-known Faraday effect—where the magnetization of a material rotates the polarization of light transmitted through it. They found that the switching process completes within a time well below 90 ps. The reversal is initiated in a small region in which the heating essentially destroys the magnetization for an instant. Subsequently, after about 30 ps, the magnetization nucleates and rebuilds as the temperature of the electrons decreases. This either leads to an expansion of the nucleus to form a larger area with a reversed magnetization or to the restoration of the initial state, depending on the initial magnetization direction and the chirality of the laser pulse.

2.6 Outline of the thesis

The magnetization switching achieved through thermal activation, stress-induced anisotropy and by applying a high intense laser pulsed field is having its own demerits like loss of information due to heating up of atoms. In the conventional magnetic recording, the reversing field is applied antiparallel to the direction of the magnetization that limits the reversal speed. The development in magnetic data storage devices, the need of increasing the data storage capacity, the speed of writing and retrieving data in memory devices lead a new revolution in scientific industries.

Ferromagnetic systems with different magnetic interactions such as bilinear, biquadratic, octupole-dipole, Dzyaloshinskii-Moriya, site-dependent interactions and interactions with the magnetic field component of the electromagnetic field have acted as important nonlinear dynamical models which introduce a special type of elementary spin excitations. The elementary spin excitations in these systems are governed by spatially confined entity called "Solitons" that is localized, stable moving objects that scatters elastically. In this direction, I intended to write this thesis to utilize the coherent and localized nature of magnetic solitonic state to reverse it without losing energy during re-
versal process. The switching of magnetization of the medium through flipping of soliton furnishes the possibility of developing new innovative approach for data storage technologies. It is also theoretically proved that inhomogeneous exchange interaction is a good candidate for inducing magnetization reversal or spin reversal process without applying any external magnetic field in the ferromagnetic systems. Therefore it has become theoretically more important in the recent years to understand the magnetization reversal process by solving the associated dynamical equations.

In the third chapter, we investigate the magnetization reversal through soliton in a site-dependent Heisenberg ferromagnetic spin chain with Gilbert damping in the presence of localized inhomogeneity in the classical continuum limit and its associated dynamics is governed by Landau-Lifshitz-Gilbert (LLG) equation. As it was found to be difficult to solve the LLG equation in its natural vector form, we tried to rewrite it in an equivalent representation by mapping the continuous spin chain onto a moving helical space curve and recast the spin equation equivalently as the evolution equations for the curvature and torsion of the space curve. More interestingly it was found that the spin dynamics in this representation is governed by generalized higher order inhomogeneous nonlinear integro-differential equation in the form of Schrödinger equation. In order to study the effect of localized inhomogeneity on the soliton, a multiple scale perturbation analysis was carried out and obtained a system of coupled evolution equations for the velocity and amplitude. We employed Jacobi elliptic function method to solve the evolution equations for soliton parameters. The evolution of the amplitude and velocity of the soliton leads to magnetization reversal via flipping of solitons in the ferromagnetic medium and this switching behaviour of soliton is also demonstrated through fourth order Runge-Kutta method numerically. Finally, we have studied the nature of perturbed soliton solutions.

The energy-momentum transport phenomenon through soliton along the inhomogeneous one-dimensional ferromagnetic spin chain with relativistic Gilbert damping for linear inhomogeneity is presented as second section of this chap-
The influence of inhomogeneity and damping on the evolution of energy and current densities of the magnetization is demonstrated. It is found that the presence of inhomogeneity and damping support the loss-less energy-momentum transport along the site-dependent spin chain.

The nonlinear dynamics of a site-dependent Heisenberg one-dimensional ferromagnetic spin chain with Gilbert damping is expressed in the form of Landau-Lifshitz-Gilbert equation in the classical continuum limit. We stereographically project the unit sphere of spin onto a complex plane and the associated underlying nonlinear spin dynamics can be expressed in terms of new variable. From the evolution equation in terms of stereographic variable, the exact spin soliton solution is constructed with the aid of symbolic computation using modified extended tanh-function method. The effect of inhomogeneity and damping on the spin soliton is studied and it is presented as third section in the third chapter.

In the fourth chapter, we investigate the magnetization reversal through soliton in a one-dimensional crystal field anisotropic Heisenberg ferromagnetic spin chain with varying bilinear and biquadratic exchange interactions in the semi-classical limit. The associated spin dynamics are governed by an inhomogeneous higher order nonlinear Schrödinger equations in the semi-classical limit through Glaubers coherent state method combined with the Holstein-Primakoff bosonic representation for spin operators and the integrability condition is also established through Painlevé test. A direct multiple scale perturbation analysis is carried out for generalized higher order nonlinear Schrödinger equation and evolution equations for soliton parameters is obtained in terms of velocity and amplitude. From the solutions, it is found that the velocity and amplitude of soliton exhibits controlled magnetization reversal behavior in the nanoscale regime. This switching behaviour of soliton has been discussed in this chapter.

In this chapter, we study the nonlinear dynamics of an inhomogeneous bilinear and biquadratic isotropic Heisenberg ferromagnetic spin chain in the classical continuum limit is governed by Landau-Lifshitz (LL) equation. This
dynamical equation is equivalently represented by mapping it onto a moving helical space curve and is governed by a higher order generalized NLS equation. The demonstration of the creation and annihilation of soliton through the sine-cosine function method under the influence of nonlinear inhomogeneities on the spin chain is also discussed.

The combination of low symmetry and spin-orbit coupling was shown by Dzyaloshinskii and Moriya to give rise to anisotropic exchange coupling leading to the mechanism of weak ferromagnetism results in the canting of spins. Dzyaloshinskii has shown that the spin superstructure gives rise to a non-vanishing antisymmetric spin coupling vector which is parallel to the trigonal axis in $\alpha - Fe_2O_3$. Moriya has shown how the process involving an additional virtual transition due to spin-orbit coupling can cause an anisotropic exchange interaction as a correction to the isotropic Anderson superexchange term and introduced a new term in the spin Hamiltonian which is the Dzyaloshinskii-Moriya interaction. The fifth chapter presents the magnetization reversal process through soliton flip under the influence of an antisymmetric Dzyaloshinskii-Moriya interaction in an anisotropic one-dimensional site-dependent ferromagnetic spin chain in the semi-classical limit which is governed by a generalized inhomogeneous nonlinear Schrödinger equation in the continuum limit. Then the evolution soliton parameters in terms of amplitude and velocity has been obtained through the nonlinear multiple scale perturbation analysis. We demonstrate the magnetization reversal through flipping of soliton in the ferromagnetic medium by solving the two coupled evolution equations for the velocity and amplitude of the soliton using the fourth order Runge-Kutta method numerically. We propose a novel approach to control the flipping behaviour of soliton by suitably tuning the parameter associated with Dzyaloshinskii-Moriya interaction which forces the soliton to move with constant velocity and amplitude along the spin lattice. Finally, the perturbed one soliton solution has been constructed in this chapter.

The sixth chapter presents the study of magnetization reversal through flipping of soliton in a one-dimensional anisotropic Heisenberg ferromagnetic spin
chain with octupole-dipole interaction. The nonlinear dynamics associated with octupole-dipole interaction is governed by a generalized higher order nonlinear Schrödinger (GNLS) equation obtained through semi-classical approach in the continuum limit. In order to understand the underlying nonlinear dynamics in the form of solitons, we solve the GNLS equation through multiple scale perturbation analysis. Finally, we have investigated the nature of evolution of perturbed soliton solutions.

The magnetization reversal through soliton in a one-dimensional anisotropic site-dependent bilinear ferromagnet under the influence of electromagnetic field in the classical continuum limit has been discussed in the seventh chapter. It is found that the associated spin dynamics coupled with Maxwell equations are governed by an inhomogeneous nonlinear Schrödinger equation through reductive perturbation method in which the components of the magnetization of the medium and that of applied electromagnetic wave are perturbed in a nonuniform way. In order to study the effect of inhomogeneity on the spin soliton under the influence of EM wave, the multiple scale perturbation analysis is carried out. From the analysis, it is found that the velocity of the soliton undergoes magnetization reversal behavior in the nano-scale regime due to the presence of linear inhomogeneity whereas the amplitude of soliton remains constant during the course of time. Finally, the magnetization of the medium and hence the magnetic field components of EM wave were constructed and its nature of evolution is discussed in this chapter.

The propagation of electromagnetic soliton in a nonlinear anisotropic biquadratic ferromagnetic medium which are free from electric charges is studied in the seventh chapter. It is found that the dynamics of the system was obtained by solving the Landue-Lifshitz equation coupled with Maxwell equations through reductive perturbation method. We employed the multiple scale perturbation on the higher order nonlinear Schrödinger equation and found that the velocity and amplitude of the soliton remains constant when passes through the biquadratic ferromagnetic medium. The associated components of magnetization of the medium is constructed and its nature of evolution has
been discussed in this chapter.

In the eighth chapter, we investigate the nonlinear dynamics of an isotropic spin ladder with two ferromagnetic legs coupled antiferromagnetically and the associated spin dynamics are governed by a set of two coupled Landau-Lifshitz equations. Then we rewrite the two coupled LL equations in terms of stereographic variables by projecting the unit sphere of spins onto two complex planes. The evolution equations associated with LL equation is solved through modified tangent hyperbolic function method. From the evolution of spin components, we demonstrate sharing of energy among the legs via rung and thus soliton exhibits switching between bistable states leading to the possibility of construction of logic gates.