CHAPTER ONE

Time-Domain Study of Magnetization Dynamics in Magnetic Thin Films and Micro- and Nanostructures

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1. GENERAL INTRODUCTION

Magnetism is an old subject whose origin goes back to about 600 BC [1]. Since its early discovery, magnetism has been used to the benefit of human society starting from its use in surgery to the use of compass for navigation. Over the years, magnetism has seen discovery of many fundamental phenomena and contributed many modern-day technologies in electromagnetics industry, health science, and more recently in miniaturized solid-state devices including magnetic data storage and memory devices. The modern
magnetism has its origin in the beginning of nineteenth century, which began with the seminal works by Oersted, Ampere, Gauss, Biot-Savart, Faraday, and Lorentz, which formed the basis of magnetostatics and magnetodynamics. Based upon these works, Maxwell [2] in 1861–1862 formulated the well-known Maxwell’s equations unifying electricity, magnetism, and optics.

Magnetic domain theory was developed in 1906 by Weiss [3], who suggested that a large number of atomic magnetic moments are aligned parallel in each domain. Weiss explained the spontaneous alignment of atomic moment within a ferromagnetic material by Weiss’s mean field theory. If a system is in a single-domain state in equilibrium and the spins rotate uniformly during reversal, then it can be described by a macrospin model. Quasisingle-domain systems may reverse by inhomogeneous modes, such as curling, fanning, buckling, and by the formation of “C,” “S,” and vortex states [4,5]. The simplest classical model describing single-domain systems with coherent rotation was developed by Stoner and Wohlfarth [6] and is commonly referred to as Stoner–Wohlfarth model. The aim of this model is to analytically calculate the equilibrium directions of the magnetization for a given anisotropy and a given applied field and its history. Further, for a magnetic field applied along a given direction, one can calculate the field value at which the magnetization reverses. As spatial variation of magnetization is ignored in this model, the exchange energy does not play a role and the magnetic switching is governed by the interplay between the Zeeman energy and the effective anisotropy. In 1935, Landau and Lifshitz [7] gave a theory of the dispersion of magnetic permeability in ferromagnetic body and predicted an equation of motion describing precession of magnetization. The Landau–Lifshitz equation was later modified by Gilbert [8], where the original damping term was replaced by a dimensionless Gilbert damping term. In 1946, Griffith [9] experimentally observed the ferromagnetic resonance (FMR) for the first time. In 1948, Kittel [10] derived the formula for the FMR frequency in terms of external magnetic field and the internal magnetic parameters famously known as Kittel formula.

Today, with the invention of various new magnetic materials, synthetic structures, micro- and nanostructures, and metamaterials, magnetism has come a long way and found applications in a range of multidisciplinary fields in modern and future nanotechnologies like nonvolatile magnetic memory [11,12], magnetic storage media [13], magnetic recording heads [14], magnetic resonance imaging [15], and biomedicine and health science [16,17]. More recently, emerging technologies are proposed in the fields of spin logic [18,19], spin-torque nano-oscillators (STNOs) [20], and magnonic
crystals [21]. The new technologies demand the invention of new material properties, which may not always be possible to achieve in a single material in the bulk or thin-film form. Instead, structuring of known materials in one and two dimensions at various length scales and exploiting dynamical magnetic properties over various time scales may potentially offer the desirable material properties. An existing example is the patterned magnetic media, which uses ordered arrays of lithographically patterned two-dimensional arrays of bits, and the magnetic switching behaviors of such systems including switching field distribution have been thoroughly studied. An essential criterion has been to eliminate magnetostatic interaction (crosstalks) between the individual bits for the application in patterned magnetic media. The same system may be used to transmit magnetic excitations in the form of collective long wavelength spin waves as information carrier in magnonic crystals, when the individual magnetic elements are strongly magnetostatically coupled. For various applications, exploration of a variety of new phenomena is required and this ranges from slower processes such as domain wall dynamics and magnetic vortex dynamics to faster processes such as spin-wave propagation and localization and ultrafast demagnetization and relaxation. This also introduces magnetic structures at various length scale such as nanodots, microdisks, and finally magnetic nanowires and nanostripes.

The key challenges to investigate and to apply the physical phenomenon mentioned above are the synthesis or fabrication of high quality magnetic materials and their characterization. Technology demands fabrication of magnetic ultrathin films and multilayers (MLs) with high surface and interface qualities. Applications in various magnetoelectronic and magnonic devices requires the fabrication of nanomagnets with narrow size dispersion arranged in an ordered array over a macroscopic length scale. These led to the development of a number of “top-down” and “bottom-up” approaches in nanofabrication methods and more recently a combination of the two. While the bottom-up approach mainly relies upon solution phase colloidal chemistry [22] and electrochemistry using different templates such as track-etched polymer [23], anodic alumina [24], and diblock copolymer membranes [25], the top-down approach relies primarily on physical processes. This includes different kinds of lithographic techniques such as photolithography [26], electron beam lithography (EBL) [27], deep ultraviolet lithography (DUV) [28], X-ray lithography [29], interference or holographic lithography (IL) [30], nanoimprint lithography (NIL) [31], and ion beam lithography (IBL) [32]. Scanning probe lithography [33], step growth methods [34], shadow masks [35], and laser or ion irradiation [36] are also very promising techniques.
The quasistatic and ultrafast magnetization dynamics of magnetic thin films, multilayers, and nanostructures are different from their bulk counterparts. Magnetization dynamics of these systems strongly depends upon their static magnetization states, which depend not only on their intrinsic material parameters such as exchange stiffness constant, saturation magnetization, and magnetocrystalline anisotropy, but also on their physical structures as well as the external parameters such as the strength and direction of the bias magnetic field. To study the quasistatic and ultrafast dynamic properties of nanomagnets, different kinds of sensitive characterization techniques have been developed in last few decades [37]. Magnetic force microscopy (MFM) [38] and Lorentz force microscopy [39] are two such examples, which are now extensively used to map the gradient of the stray magnetic field and the sample magnetization, respectively, with a spatial resolution better than 10 nm. For MFM, the contrast of the images comes from magnetic force between the scanning magnetic tip and the gradient of the stray magnetic from the sample. On the other hand, the deflection of the accelerated electrons by Lorentz force after transmitting through thin magnetic specimens creates the magnetic contrast in case of Lorentz microscopy. However, it is not straightforward to extract quantitative information directly from either of these imaging techniques. Electron holography [40] is another imaging technique based upon the electron interference, by which the amplitude and phase information of the spin configurations and stray magnetic fields can be mapped with a very high spatial resolution down to 2 nm. Magneto-optical Kerr effect (MOKE) microscopy [41] is an old technique and is widely used to map the sample magnetization with a sub-μm spatial resolution. The photoemission electron microscopy (PEEM) [42] is a form of X-ray microscopy and has a far better spatial resolution than visible light imaging. Spin–polarized low-energy electron microscopy (SPLEEM) [43], scanning electron microscopy with polarization analysis (SEMPA) [44], spin–polarized scanning tunneling microscopy (SP–STM) [45], and ballistic electron magnetic microscopy (BEMM) [46] are other imaging techniques, which give excellent spatial resolution of 10 nm or better. These techniques use the spin–dependent transmission, scattering or tunneling of electrons for image contrast. However, despite having very good spatial resolution, all the above techniques suffer from poor or moderate temporal resolution.

Subsequently, different kinds of experimental techniques have emerged to investigate fast magnetization dynamics of magnetic thin films and nanostructures. The conventional FMR [47] and vector network analyzer-based broadband ferromagnetic resonance (VNA–FMR) [48]
techniques are very efficient to measure the high-frequency magnetic resonance, permeability, and loss from MHz to tens of GHz regime limited only by the instrumental bandwidth with a very good spectral resolution. Later, Tamaru et al. [49] developed the spatially resolved FMR to image resonant mode profiles of confined magnetic elements. Pulsed inductive microwave magnetometry (PIMM) [50] is an oscilloscope-based time-domain technique to measure the magnetization dynamics with tens of picosecond temporal resolution. Brillouin light scattering (BLS) is based upon the inelastic light scattering from spin waves and other quasiparticles and is conventionally used to study the spin waves in the wave vector domain [51]. The frequency dispersion of the spin waves with their wave vector can be measured directly by varying the angle of incidence of light with respect to the sample plane using this technique. Recently space-resolved and time-resolved BLS techniques have been developed to obtain sub-\(\mu\)m spatial resolution and few ns temporal resolution [52]. The best spatiotemporal resolution is obtained from time-resolved magneto-optical Kerr effect (TRMOKE) microscope [53]. They are used to probe the ultrafast magnetization dynamics in time domain and can achieve tens of femtosecond temporal resolution limited only by the pulse width of the laser. The magnetoresistive methods [54] and X-ray microscopy [55] also have the potential to achieve very good spatiotemporal resolution similar to TRMOKE. Time-resolved scanning Kerr microscopy (TRSKM) is a variant of TRMOKE, which is used to image the time evolution of magnetization excited by a time-dependent magnetic field.

Here, we will review the time-domain study of magnetization dynamics in magnetic thin films and nanostructures. The remaining part of the review will be divided into 11 subsections. In Section 2, we will describe the background theory of magnetization dynamics. In Section 3, we will revisit the development of TRMOKE experiments from early 1990s to cater the need of various measurement systems and properties starting from fundamental science to industrial needs. We will discuss some milestone instrumental development and investigation of new phenomena using this technique. In Sections 4–9, we will describe the time-domain measurements of precessional dynamics in some specific and important magnetic systems. This includes microscopic elements (Section 4), magnetic multilayers (Section 5), magnetic nanodot arrays (Section 6), single nanomagnet (Section 7), and magnetic vortex and domain wall dynamics (Section 8). In Section 9 we review the coherent suppression of magnetization dynamics using magnetic field pulse shaping, and in Section 10 we review the progress
in precessional switching. Finally, we conclude and discuss the future directions in the study of time-domain magnetization dynamics (Section 11).

2. THEORETICAL BACKGROUND

2.1. Time Scales of Magnetization Dynamics

Magnetization dynamics can occur over a wide range of time scale. Figure 1.1 shows various kinds of magnetization dynamics with their characteristic time scales [56]. The time scales (τ) are determined by the interaction energies (E) via Heisenberg relation \( \tau = \hbar / E \). The fastest process is the fundamental exchange interaction, which occurs within 10 fs. The spin–orbit coupling and spin-transfer torque occur in the time scale of 10 fs–1 ps. Laser-induced ultrafast demagnetization occurs within few hundreds of fs. The fast remagnetization time following the ultrafast demagnetization covers the time scale of 1–10 ps. The magnetic writing which is done via reversal of spin has a time scale of few ps to few hundreds of ps, whereas vortex core switching occurs from few tens of ps to ns time scale. The precession of magnetization occurs within few ps to few hundreds of ps whereas the damping associated with magnetization precession occurs from sub-ns to tens

![Figure 1.1 Characteristic times scales of various kinds of magnetization dynamics.](image-url)
of ns time scale. The spin waves in ferromagnetic material can propagate in a time scale of few hundreds of ps to tens of ns time scale before it dies out. The slowest process is the domain wall motion, which has the time scale from few ns to hundreds μs.

2.2. Laser-Induced Ultrafast Magnetization Dynamics

The researchers have been attracted to the question: what happens when an ultrashort laser pulse interacts with an ordered ferromagnetic material. The breakthrough in this research topics occurred after the pioneering work by Beaurepaire et al. [57] in 1996. They found that a nickel thin film can be demagnetized in a subpicosecond time scale after excitation with a sub-100fs pulsed-laser beam. A number of studies on all elementary ferromagnetic transition metals (Co, Ni, Fe) and several alloys thereof have confirmed the above result [58]. Now the question is how the magnetic moment of a system can be quenched so rapidly, while its total angular momentum is conserved. Before laser excitation, the angular momentum is mainly carried by the aligned electron spins. Thus, a transfer of angular momentum to another degree of freedom, such as electrons or lattice, is required to compensate the loss of magnetic order [59]. Figure 1.2 shows a sequence of phenomena observed when a ferromagnetic material is irradiated with ultrashort laser pulses. The underlying mechanism is discussed by considering the transfer of angular momentum among different degrees of freedom.

Figure 1.2 The schematic diagram of sequence of mechanisms in ultrafast magnetization dynamics.
A number of processes occur when a femtosecond laser pulse interacts with an ordered ferromagnetic material. First of all, the photon field of laser beam interacts with spin degrees of freedom of electrons. The angular momentum of light is modified nonlinearly within first 50 fs. The electronic distribution is not initially thermalized. The electrons are excited above the Fermi level because of electron–electron scattering [60]. The excited electrons are called “hot electrons” [61]. Within this time the electronic temperature may increase up to $10^3$ K [60]. Shortly after that the thermalization of spin population occurs as the hot electrons excite the spins by electron–magnon interaction [62,63]. Consequently, there is a difference between the thermalization times of charges ($\tau_e$) and spins ($\tau_s$). During the thermalization of charges and spins, the electronic wavefunction loses its phase memory (incoherent process) with respect to the excitation, which leads to the demagnetization of the ferromagnet. However, the underlying mechanism of demagnetization is still a subject of intense debate.

A majority of the magnetism community believe that during the excitation and thermalization of the electrons, Stoner pairs with a large kinetic energy are excited in ferromagnetic metals like nickel or cobalt [60]. Therefore, most of the published reports on the ultrafast demagnetization state that ultrafast spin–flip scattering plays important role for demagnetization. The primary one is the Elliott–Yafet type of scattering, which says that an electron flips its spin due to the influence of impurity centers and phonons and emits or absorbs a phonon [59,64]. Other kinds of scattering processes are electron–magnon spin–flip scattering [65], Coulomb exchange spin–flip scattering [66], and relativistic electromagnetic radiation-induced spin–flip scattering [67]. However, few reports claim that other kinds of physical mechanisms like excitation of Stoner pairs [68], the spin–orbit coupling [69], and the coupling with the electromagnetic field via a terahertz emission [70] are responsible for ultrafast demagnetization. Moreover, some recent reports show that the superdiffusive spin transport rather than spin–flips plays the major role in the magnetization processes on the femtosecond time scale [71,72].

Following the demagnetization, the electronic charges and spins start to relax. This relaxation occurs in two different time scales (few ps to ns). Therefore, the reflectivity (depends on the population of charges above Fermi level) and also magnetization of the system decay biexponentially. The faster relaxation time ($\tau_1$) of the hot electrons and spins occurs, because they exchange energy with the lattice and this is mediated by the electron–phonon interaction. This fast relaxation time ($\tau_1$) may vary from few ps to
tens of picosecond depending on the electron–phonon coupling and also the specific heat of the electrons and phonons \[60,73,74\]. The relaxation time also depends on the density of laser excitation and magnetocrystalline anisotropy of the material to some extent. In this time scale the thermodynamic quantities such as specific heat and temperature come into play. Therefore, one may talk about phenomenological three temperature model (3TM) \[57,75\]. In 3TM, heat capacities and temperatures are assigned to the reservoirs of electronic charge (e), spin (s), and lattice/phonons (p): \((C_e; T_e)\), \((C_s; T_s)\), and \((C_p; T_p)\), respectively. Furthermore, coupling constants between charge–spin, spin–lattice, and charge–lattice are defined as \(g_{cs}\), \(g_{sp}\), and \(g_{cp}\). They describe the rate of energy exchange between the participating subsystems. Thus, the overall dynamics is phenomenologically described by a set of three coupled differential equations (rate equations). By exchanging heat, spins come to an equilibrium temperature with charge and lattice. Therefore, the lattice temperature changes. The magnetocrystalline anisotropy, which is a function of lattice temperature, also changes. The ultrafast change in magnetocrystalline anisotropy acts as an effective pulsed field and triggers the precession of magnetization. The second or longer relaxation time \(\tau_2\) corresponds to the diffusion of electron and lattice heat to the surroundings (such as substrate). In this time scale the precession of magnetization also damps out. These mechanisms can be described by Landau–Lifshitz–Gilbert (LLG) formalism.

2.3. LLG Equation

The dynamics of a spin can be described mathematically by an equation of motion. The equation of motion can be derived from quantum mechanics \[76,77\]. The equation of motion for a single spin can be written as

\[
\frac{d}{dt}\langle S \rangle = \frac{g\mu_B}{\hbar} (S \times B) \tag{1.1}
\]

In the macrospin model the magnetization \(M\) is supposed to be uniform throughout the sample. The relation between \(S\) and \(M\) can be written as

\[
M = -\frac{g\mu_B}{\hbar} \langle S \rangle \tag{1.2}
\]

Therefore, equation of motion of magnetization in presence of external magnetic field will be
\[
\frac{dM}{dt} = -\frac{g\mu_B}{\hbar} (M \times H) = -\gamma (M \times H),
\]  

(1.3)

where \( \gamma = g\mu_B/\hbar \), the gyromagnetic ratio. The above equation is called the Landau–Lifshitz equation (without damping).

The above equation can be generalized by replacing \( H \) by \( H_{\text{eff}} \). This equation implies that the tip of the magnetization vector precesses around the effective magnetic field in a circular orbit as shown in Fig. 1.3a for infinitely long time with an angular frequency \( \omega = \gamma H_{\text{eff}} \). Practically, the precession amplitude of magnetization decreases with time and the tip of the magnetization vector follows a spiral path (Fig. 1.3b). Therefore, a damping or relaxation term should be added with the Landau–Lifshitz equation. Landau–Lifshitz suggested a damping term as \((\lambda/M_s^2)M \times (M \times H_{\text{eff}})\). Gilbert suggested a damping term as \((a/M_s)(M \times (dM/dt))\).

With the above expression, the equation of motion of magnetization vector can be expressed by the following equation:

\[
\frac{dM}{dt} = -\gamma (M \times H_{\text{eff}}) + \frac{a}{M_s} \left( M \times \frac{dM}{dt} \right).
\]

(1.4)

Equation (1.4) is known as LLG equation. Here \( a \) is the Gilbert damping parameter. The value of \( a \) is much less than 1 for transition metals, which ensures that the magnetization precesses a number of times around the effective field before coming to the equilibrium position.

### 2.4. Ferromagnetic Resonance

In the above section we have seen that when a steady magnetic field is applied to a ferromagnetic material, the magnetization starts to precess around the effective magnetic field with an angular frequency \( \omega = \gamma H_{\text{eff}} \).

![Figure 1.3 Precession of magnetization vector \( M \) around magnetic field \( H \) (a) in absence of damping term and (b) in presence of damping term.](image-url)
If an alternating magnetic field with same angular frequency ($\omega$) is applied in transverse direction to the steady field, then resonance will occur and the magnetization will precess in the resonance frequency, absorbing power from the alternating field. This phenomenon is called ferromagnetic resonance (FMR).

### 2.4.1 Macrospin Model of FMR: Kittel Formula

In macrospin model, the magnetization of a ferromagnetic element is considered to be uniform throughout the element. Therefore, the magnetic moment of the whole element may be represented by a giant magnetic moment (macrospin). In this situation the resonance condition of a ferromagnetic element can be derived theoretically. Griffith [9] observed that the resonance frequencies were two to six times greater than the calculated Larmor frequencies for electron spin. Later, Kittel [78] explained that the dynamic coupling caused by the demagnetizing field normal to the specimen surface should be taken into account. Consider a ferromagnetic specimen with a plane surface. A steady bias field $H_z$ is applied along the $z$-axis and rf field $H_x$ is applied along the $x$-axis. The equation of motion for this system using Eq. (1.3) can be written as

$$
\frac{\partial M_x}{\partial t} = \gamma (H_z + 4\pi M_z) M_y = \gamma B_z M_y,
$$

$$
\frac{\partial M_y}{\partial t} = \gamma (M_z H_x - M_x H_z),
$$

$$
\frac{\partial M_y}{\partial t} \approx 0. \tag{1.5}
$$

Solving these equations by considering time-dependent variation of $M$ and $H(\exp(j\omega t))$, the resonant frequency for a plane surface comes out to be

$$
\omega_0 = \gamma (B_z H_z)^{1/2}. \tag{1.6}
$$

The resonant frequencies for other shapes will be different after taking into account the demagnetizing factors along three coordinate axes. For a specimen of ellipsoidal shape and with principal axes parallel to $x$, $y$, and $z$ axes, the effective magnetic field values inside the specimen can be written as

$$
H'_x = H_x - N_x M_x,
$$

$$
H'_y = -N_y M_y,
$$

$$
H'_z = H_z - N_z M_z. \tag{1.7}
$$
Here \( N_x, N_y, N_z \) are the demagnetizing factors. Substituting effective field values in Eq. (1.3), the components can be written as

\[
\frac{\partial M_x}{\partial t} = \gamma [H_z + (N_y - N_z)M_z]M_y, \\
\frac{\partial M_y}{\partial t} = \gamma [M_z H_x - (N_x - N_z)M_x M_y - M_x H_z], \\
\frac{\partial M_y}{\partial t} \approx 0.
\] (1.8)

For this case, the resonant frequency becomes

\[
\omega_0 = \gamma \left( [H_z + (N_y - N_z)M_z] \times [H_z + (N_x - N_z)M_z] \right)^{1/2}.
\] (1.9)

Here, magnetocrystalline or other forms of magnetic anisotropy are not considered, which may be present in the system. If anisotropy terms are expressed as effective demagnetizing factors \((N_x^e, N_y^e, N_z^e)\), then the resonant frequency can be written in the following form:

\[
\omega_0 = \gamma \left( [H_z + (N_y^e - N_z)M_z] \times [H_z + (N_x^e - N_z)M_z] \right)^{1/2}.
\] (1.10)

Equation (1.10) is known as the Kittel formula, which gives the resonant frequency of magnetization for a uniformly magnetized ferromagnetic substance.

### 2.4.2 Effective Demagnetizing Factors: Twofold and Fourfold Anisotropy

In the following, we will find out the demagnetizing factors for two- and fourfold anisotropies by considering an orthogonal system with \( a, b, \) and \( c \) axes [79]. The geometry of the magnetization \((\mathbf{M})\) and applied magnetic field \((\mathbf{H})\) is shown in Fig. 1.4.

If \( F \) represents the anisotropy energy, then we can write

\[
\nabla F = \left( \frac{\partial F}{\partial \theta} \right) \hat{\theta} + \frac{1}{\sin \theta} \left( \frac{\partial F}{\partial \phi} \right) \hat{\phi},
\] (1.11)

Torque, \( \mathbf{I} = -\hat{r} \times \nabla F = -\left( \frac{\partial F}{\partial \theta} \right) \hat{\phi} + \frac{1}{\sin \theta} \left( \frac{\partial F}{\partial \phi} \right) \hat{\theta}. \) (1.12)

If we consider a thin film in the \( a-b \) plane with an in-plane uniaxial anisotropy and with the magnetic field applied in the plane of the film \((\theta = \pi/2)\), then we may write
\[
\left( \frac{\partial F}{\partial \theta} \right)_{(\pi/2)+\alpha} - \left( \frac{\partial F}{\partial \theta} \right)_{\pi/2} = N_{x}^{c}M_{s}^{2}\sin\alpha \tag{1.13}
\]

and

\[
\left[ \left( \frac{\partial F}{\partial \theta} \right)_{\phi + \beta} - \left( \frac{\partial F}{\partial \theta} \right)_{\phi} \right]_{\theta = \pi/2} = N_{y}^{c}M_{s}^{2}\sin\beta. \tag{1.14}
\]

If the anisotropy energy is considered to be along \(a\)-axis, then it can be written as \(F_{2} = -K_{2}\sin\theta\cos^{2}\phi\) and the resonant frequency can be calculated to be

\[
\omega_{0} = \gamma \left[ H_{z} + \left( N_{y} - N_{z} + \frac{2K_{2}}{M_{s}}\cos 2\phi \right) M_{z} \right] \times \left[ H_{z} + \left( N_{x} - N_{z} + \frac{2K_{2}}{M_{s}}\cos 2\phi \right) M_{z} \right]^{1/2}. \tag{1.15}
\]

For fourfold anisotropy, the free energy can be written as \(F_{4} = (K_{4}/4)\sin^{4}\theta(3 + \cos 4\phi)\) and the resonant frequency can be calculated as

\[
\omega_{0} = \gamma \left[ H_{z} + \left( N_{y} - N_{z} - \frac{4K_{4}}{M_{s}}\cos 4\phi \right) M_{z} \right] \times \left[ H_{z} + \left( N_{x} - N_{z} - \frac{K_{4}}{M_{s}}(3 + \cos 4\phi) \right) M_{z} \right]^{1/2}. \tag{1.16}
\]
2.5. Spin waves

2.5.1 Introduction
The concept of spin wave was first introduced by Bloch [80] in 1930. If an array of exchange- and dipolar-coupled spins is disturbed locally by some external perturbation, then the disturbance propagates as a wave generated by collective phase coherent precession of spins in the ferromagnetic medium. The quanta of spin waves are called magnons. Classically magnetization dynamics is governed by the Landau–Lifshitz equation (neglecting damping), which shows that the rate of change of magnetization is proportional to the torque exerted by the effective magnetic field ($H_{\text{eff}}$) on the magnetization. Both dipolar and exchange interactions contribute to the effective magnetic field.

2.5.2 Exchange Spin Wave: Dispersion Relation

The spin wave is dominated by exchange interaction when the wavelength ($\lambda$) of spin wave becomes very short (i.e., the wave vector ($k$) becomes very large). In the ground state all the spins are parallel to each other. For a chain of $N$ spins with magnitude “$S$” for each of them, the Heisenberg exchange interaction energy of $p$th spin will be

$$E_p = -2JS_p \cdot (S_{p-1} + S_{p+1})$$

The torque on the $p$th spin can be written as

$$\frac{dS_p}{dt} = \left(\frac{2J}{\hbar}\right) (S_p \times S_{p-1} + S_p \times S_{p+1}).$$

Considering harmonic solutions, that is,

$$S_p^x = me^{i[pka-\omega t]}, \quad S_p^y = me^{i[pka-\omega t]},$$

the dispersion relation can be obtained as

$$\hbar \omega = 4JS(1 - \cos ka),$$

where $k$ is the wave vector and $a$ is the lattice constant. At long wavelength limit, $ka \ll 1$ and Eq. (1.20) can be written as

$$\hbar \omega = (2JSa^2)k^2.$$
The dispersion relation is quadratic in nature for long wavelength limit, that is, in small $k$ limit. The dispersion of exchange spin wave is isotropic in nature.

### 2.5.3 Exchange Spin Waves in Thin Films

Perpendicular standing spin-wave (PSSW) mode observed in a thin film is an example of exchange spin wave. The spin wave propagates along the direction perpendicular to the plane of film and reflected back to interfere and form standing spin wave. The wave vector is quantized for pinned or unpinned boundary condition. The wave vector can have the value $n(\pi/d)$, where $d$ is the film thickness and $n$ is the positive integer. In Fig. 1.5a, the pattern of PSSW mode is shown for $n = 0$ and $n = 1$.

### 2.5.4 Magnetostatic Modes in Thin Films

If a magnetic field is applied in the plane of an infinite ferromagnetic thin film, the magnetic moments are all aligned in the plane of the film. The wavelength of the excited spin waves depends upon the characteristic length scale of uniformity of the internal magnetic field. As the magnetization and hence the internal magnetic field for an infinite thin film is uniform, therefore large wavelength spin waves are observed for such thin films. The spin waves observed in this case are governed by dipolar interaction. The dispersion relations of dipolar modes can be calculated numerically by solving

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**Figure 1.5** (a) Schematics of perpendicular standing spin-wave and magnetostatic surface wave modes are shown for a ferromagnetic thin film. (b) The calculated dispersions of three different magnetostatic spin-wave modes are shown for a permalloy (Ni$_{80}$Fe$_{20}$) thin film with $4\pi M_s$ value 10.8 kOe. The value of applied bias was chosen as 1 kOe for DE and BWVMS, and 11 kOe for FWVMS in the calculation.
Landau–Lifshitz equation after considering Maxwell’s equations in magnetostatic limit:
\[
\nabla \times \mathbf{H} = 0,
\quad \nabla \cdot (\mathbf{H} + 4\pi \mathbf{M}_s) = 0.
\]

(1.23)

The magnetization can be written as
\[
\mathbf{M}(\mathbf{r}, t) = \mathbf{M}_s + \mathbf{m}(\mathbf{r}, t).
\]

(1.24)

Here, \( \mathbf{M}_s \) is the saturation magnetization and \( \mathbf{m} \) is a small variation of magnetization due to the precession of magnetization. For small angle precession \( \mathbf{m} \ll \mathbf{M}_s \) and \( \mathbf{m}(\mathbf{r}, t) \) can be written as sum of a series of plane waves:

\[
\mathbf{m}(\mathbf{r}, t) = \sum_k m_k e^{i(\omega t - \mathbf{k} \cdot \mathbf{r})}.
\]

(1.25)

Under this condition the nonlinear Landau–Lifshitz equation can be linearized and solved. Herrings and Kittel derived the dispersion relation of dipole-exchange spin wave in an infinite ferromagnetic material as given by the formula below [81]:

\[
\omega = \gamma \left[ \left( \frac{H}{M_s} + 2A k^2 \right) \left( H + \frac{2A}{M_s} k^2 + 4\pi M_s \sin^2 \theta_k \right) \right]^{1/2},
\]

(1.26)

where \( \theta_k \) is the angle between \( \mathbf{k} \) and \( \mathbf{M}_s \). For a thin film of thickness \( d \), this dispersion relation is modified due to the broken translational symmetry at the surface or interface. Considering the external field is applied along \( z \)-axis in the plane of the film (\( y-z \) plane) and \( x \)-axis is normal to the film plane, Damon and Eshbach [82] calculated the dispersions of dipolar modes for a thin magnetic film and found two types of modes namely, the surface or Damon–Eshbach (DE) mode and the volume mode. The surface modes exist up to a critical angle \( \theta_c \) with respect to the direction perpendicular to the magnetization, given by

\[
\theta_c = \tan^{-1} \sqrt{\frac{4\pi M_s}{H}}.
\]

(1.27)

The amplitude of the surface mode concentrates primarily in the vicinity of the film surface and decays exponentially away from the surface. The penetration depth \( \delta \) decreases with the increase in angle \( \theta \) and becomes zero at \( \theta = \theta_c \), when it localizes strictly at the surface. Also, \( \delta \) is inversely proportional to wave vector along that direction [83]. As \( k_y \to 0 \), the penetration
depth $\delta \to \infty$, that is, the surface mode becomes the Kittel mode. The penetration depth is independent of film thickness and therefore DE mode is converted to bulk mode when film thickness becomes comparable to the penetration depth. In general, when the surface mode propagates perpendicular to the direction of magnetization, it is called magnetostatic surface wave (MSSW) mode. Considering negligible anisotropy, the dispersion relation of MSSW is given by

$$\omega_{DE} = \gamma \left[ H(H + 4\pi M_s) + (2\pi M_s)^2 \left( 1 - e^{-2kd} \right) \right]^{1/2}.$$ (1.28)

At $k = 0$, the frequency of MSSW mode becomes identical to the uniform (Kittel) mode:

$$(\omega_{DE})_{k=0} = \gamma [H(H + 4\pi M_s)]^{1/2}.$$ (1.29)

At $k \to \infty$, the frequency of MSSW mode becomes

$$(\omega_{DE})_{k=\infty} = \gamma (H + 2\pi M_s).$$ (1.30)

There is a manifold of volume modes and each volume mode corresponds to a mode number. The volume modes are degenerate along the propagation direction perpendicular to $M$ but the degeneracy is lifted when $k$ becomes parallel to $M$. For $k \perp M$, the volume mode becomes dispersionless, while for other directions negative dispersion is observed. When $M$ and $k$ are parallel and both lie in the plane of the film, the spin wave is called backward volume magnetostatic (BWVMS) mode. Considering negligible anisotropy, the dispersion relation of the lowest order BWVMS is given by [84,85]

$$\left( \frac{\omega_B}{\gamma} \right)^2 = H \left[ H + 4\pi M_s \left( \frac{1 - e^{-2kd}}{kd} \right) \right].$$ (1.31)

The negative slope of dispersion curve (Fig. 1.5b) implies that the phase velocity and the group velocity are in opposite directions. At $k \to 0$, the frequency of BWVMS becomes same as the Kittel mode (Eq. 1.29).

When the magnetization is along the normal to the plane of the film and the propagation direction is in the plane of the film, the forward volume magnetostatic (FWVMS) modes are observed. In the long wavelength limit the dispersion relation can be expressed after neglecting anisotropy as

$$\left( \frac{\omega_F}{\gamma} \right)^2 = (H - 4\pi M_s) \left[ H - 4\pi M_s \left( \frac{1 - e^{-kd}}{kd} \right) \right].$$ (1.32)
In the magnetostatic modes, the exchange energy is generally neglected. This is a valid approximation for materials with very weak exchange stiffness constants such as yttrium iron garnet (YIG). However, for transition metal elements and alloys, the value of $A$ is much larger and hence the contribution of exchange interaction cannot be neglected. The frequency of dipole-exchange modes therefore becomes larger than the pure dipolar modes. The contribution from exchange interaction becomes greater for larger wave vector modes as exchange has a quadratic contribution to the dispersion relation. The typical dispersion relations for various magnetostatic modes are shown in Fig. 1.5b.

2.5.5 Magnetostatic Modes in Confined Thin Ferromagnetic Elements
In Section 2.5.3, we have seen that quantized PSSW mode is observed across the thickness of thin films due to the physical boundaries. The spin wave can also be quantized in the plane of the film due to the lateral confinement of ferromagnetic element. Stripes and dots are examples of such confined magnetic elements. If $w$ is the width of the element then the values of quantized wave vector spin waves may be written as [85]

$$k_n = \frac{2\pi}{\lambda_n} = \frac{n\pi}{w}. \quad (1.33)$$

However, the analysis of the quantized spin waves is not possible only by considering the allowed wave vector in the DE modes. One also has to take into account the nonuniform demagnetizing field and the edge effects [86]. The quantized spin waves can be solved easily by solving Landau–Lifshitz equation under various micromagnetic frameworks.

2.6. Magneto-optical Kerr Effect
2.6.1 Introduction
Magneto–optical effects are very old and well-known phenomena. In 1854 Michael Faraday first invented magneto–optical phenomena in a piece of glass placed in between two magnetic pole pieces. He observed that with the magnetic field applied along the propagation direction of light, the axis of polarization is rotated. The angle of rotation is a function of strength of magnetic field and distance traveled in the medium. In 1877, Kerr [87] also observed a magneto–optical effect in reflection from a polished metallic pole of an electromagnet. Both these effects originate from the interaction of light with the applied magnetic field and magnetization of the material. In 1884, Kundt [88] observed that magneto–optical effect is greatly enhanced when light is
reflected from or transmitted through a ferromagnetic material. The magneto- 
optical rotation is linearly proportional to the magnetization in the first order. 
Magneto-optical effect based techniques have proved to be very efficient and 
sensitive techniques for imaging magnetic domains and probing magnetiza-
tion dynamics at various time scales. Magneto-optical effect depends on the 
complex dielectric tensor. A plane-polarized light is converted to an ellipti-
cally polarized light due to magneto-optical effect where the major axis of 
the ellipse is rotated from the plane of polarization of the plane-polarized light.

2.6.2 Physical Origin of Magneto-Optical Effects
The origin of magneto-optical effect can be explained by Zeeman effect. Let 
us consider a system of electrons bound harmonically in an oscillating elec-
tric field and static magnetic field. The equation of motion of each electron 
can be written as

\[
m \ddot{\mathbf{r}} + b \dot{\mathbf{r}} + k \mathbf{r} = -e \mathbf{E}_0 e^{i\omega t} - e \mu_0 \mathbf{r} \times \mathbf{H},
\]

where \( b \) is the damping coefficient, \( m \) is the electron mass, and \( (k/m)^{1/2} = \omega_0 \) 
is the natural frequency of electron. The last term in this equation is the 
Lorentz force. The above equation can easily be solved. The electric dipole 
moment \( (-e \mu_0) \) is proportional to electric field \( (\mathbf{E}_0) \). The proportionality 
tensor is called polarizability tensor. For an assembly of \( N \) oscillators the 
polarizability can be expressed as \[89\]

\[
\chi_{xx} = \chi_{yy} = \frac{N e^2}{\epsilon_0 m} \frac{\omega_0^2 - \omega^2 + j \omega \gamma}{(\omega_0^2 - \omega^2 + j \omega \gamma)^2 - 4 \omega^2 \omega_L^2},
\]

\[
\chi_{xy} = -\chi_{yx} = \frac{N e^2}{\epsilon_0 m} \frac{-2 j \omega \omega_L}{(\omega_0^2 - \omega^2 + j \omega \gamma)^2 - 4 \omega^2 \omega_L^2},
\]

where \( \omega_L = e \mu_0 H / 2m \) is the Larmor precession frequency and \( \gamma = b/m \) is the 
width of resonance. \( \chi \) is a symmetric and scalar quantity in an isotropic 
medium. Therefore the off-diagonal terms of the dielectric tensor are zero.
The anisotropy in the medium may be generated from magnetic field or 
magnetization of ferromagnetic material. For a circularly polarized light 
the polarizability can be written as \[89\]

\[
\chi_{\pm} = \chi_{xx} \pm \chi_{xy} = \frac{N e^2}{\epsilon_0 m} \frac{\left( \omega_0^2 - \omega^2 + j \omega \gamma \right) \pm 2 \omega \omega_L}{(\omega_0^2 - \omega^2 + j \omega \gamma)^2 - 4 \omega^2 \omega_L^2},
\]

\[
\approx \frac{N e^2}{\epsilon_0 m} \frac{1}{\omega_0^2 - \left( \omega \pm \omega_L \right)^2 + j \omega \gamma}.
\]
The upper sign refers to right circularly polarized (RCP), whereas the lower to left circularly polarized (LCP) light. Larmor’s theorem says that if the refractive index in absence of magnetic field is \( n_0 \), then in presence of field the refractive indices can be written as

\[
\begin{align*}
n_{\pm}(\omega) &= n(\omega \pm \omega L).
\end{align*}
\] (1.37)

Therefore the refractive indices of RCP and LCP are different in presence of a magnetic field.

Let us now consider a linearly polarized light propagating along +z direction in a circularly birefringent medium. If the polarization direction is along +x axis, then the linearly polarized light can be expressed as

\[
E = \hat{e}_x E_0 e^{i(\omega t - kz)},
\] (1.38)

where \( k \) is the wave vector and \( \omega \) is the angular frequency. In a non-birefringent medium linearly polarized light can be written as a linear combination of RCP and LCP in the following way:

\[
E = \frac{E_0}{2} (\hat{e}_x + j\hat{e}_y) e^{i(\omega t - kz)} + \frac{E_0}{2} (\hat{e}_x - j\hat{e}_y) e^{i(\omega t - kz)}.
\] (1.39)

In circularly birefringent medium the above expression will be modified as

\[
E = \frac{E_0}{2} (\hat{e}_x + j\hat{e}_y) e^{i(\omega t - 2\pi n_+ z/\lambda_0)} + \frac{E_0}{2} (\hat{e}_x - j\hat{e}_y) e^{i(\omega t - 2\pi n_- z/\lambda_0)},
\] (1.40)

where \( \lambda_0 \) is the wavelength of linearly polarized light in vacuum. The above expression can be rewritten as

\[
E = \frac{E_0}{2} e^{i(\omega t - 2\pi \bar{n} z/\lambda_0)} \left[ \hat{e}_x \cos \frac{\delta}{2} + \hat{e}_y \sin \frac{\delta}{2} \right],
\] (1.41)

where \( \bar{n} = (1/2)(n_+ + n_-) \) and \( \delta = 2\pi(n_+ - n_-)z/\lambda_0 \) is the phase difference introduced between the components by the birefringence of medium. For an absorptive medium \( n_{\pm} \) is complex quantity and can be written as

\[
n_{\pm} = N_{\pm} - jK_{\pm},
\] (1.42)

where \( n_{\pm} \) and \( K_{\pm} \) both are real quantities. In this case the linearly polarized light will be converted to an elliptically polarized light where major axis of the ellipse is rotated through an angle:
\[ \theta_K = \Re \left( \frac{\delta}{2} \right) = \Re \left( \frac{\pi l}{\lambda_0} (n_+ - n_-) \right), \] (1.43)

where \( l \) is the distance of propagation.

The ellipticity is expressed as

\[ \varepsilon_K = -\tanh \left( \frac{\delta}{2} \right) = -\tanh \left( \frac{\pi l}{\lambda_0} (n_+ - n_-) \right). \] (1.44)

### 2.6.3 Phenomenological Theory

The optical properties of material can be given by Maxwell’s equations and the consecutive relations which describe the specific material properties. We need two equations:

\[ \nabla \times E = -\mu_0 \frac{\partial H}{\partial t}, \] (1.45)

\[ \nabla \times H = \varepsilon_0 \frac{\partial E}{\partial t} + J, \] (1.46)

where \( E, H, J, \mu_0, \) and \( \varepsilon_0 \) are electric field, magnetic field, electrical current density, permeability, and permittivity, respectively. The current density can be expressed as

\[ J = \varepsilon_0 \frac{\partial P}{\partial t} = \chi \cdot \frac{\partial E}{\partial t}, \] (1.47)

where \( \chi \) is the polarizability. Again dielectric constant \( (\kappa) \) and polarizabiity \( (\chi) \) are related by the relation

\[ \kappa = 1 + \chi, \] (1.48)

where \( \kappa \) and \( \chi \) are the tensorial quantities and depend on the applied magnetic field and magnetization of the material. Their nonscalar nature gives rise to the magneto-optical effect. Now we consider a plane wave propagating through a medium. The electric and magnetic field may be written as

\[ E = E_0 e^{i(\omega t - k \cdot r)}, \quad H = H_0 e^{i(\omega t - k \cdot r)}. \] (1.49)

Substituting Eq. (1.49) in Eqs. (1.45) and (1.46), we get

\[ \left\{ n^2 1 - \kappa k / k_0^2 - \kappa \right\} \cdot E_0 = 0. \] (1.50)
Equation (1.50) possesses a nontrivial solution for $E_0$ only if the determinant of coefficients vanishes. Let us consider a cubic crystal with dc magnetic field applied along $z$-axis. Then, the dielectric constant can be written as

$$\kappa = \begin{bmatrix} \kappa_1 & \kappa_2 & 0 \\ -\kappa_2 & \kappa_1 & 0 \\ 0 & 0 & \kappa_3 \end{bmatrix}. \quad (1.51)$$

In the above equation, $\kappa_2$ is the odd function of $H$ or $M$, whereas $\kappa_1$ and $\kappa_3$ are even functions. With the above form of $\kappa$, the secular determinant for 1.50 becomes

$$n^4 \left\{ \kappa_1 + (\kappa_3 - \kappa_1) \cos^2 \theta' \right\} - n^2 \left\{ (\kappa_1^2 + \kappa_2^2 + \kappa_1 \kappa_3) - (\kappa_1^2 + \kappa_2^2 - \kappa_1 \kappa_3) \cos \theta' \right\} + \kappa_3 (\kappa_1^2 + \kappa_2^2) = 0,$$  

where

$$\cos \theta' = k_z/k = k_z/nk_0. \quad (1.53)$$

Let us now consider a simplest case: a homogeneous wave (real and imaginary parts of $k$ are parallel) with propagation direction along $H$ or $M$. In this case $\cos \theta' = 1$. Therefore the solutions of Eq. (1.53) are

$$n_{\pm} = \kappa_1 \pm j \kappa_2. \quad (1.54)$$

For simplicity, we now consider the polar Kerr effect at normal incidence. Let $k_0$, $k_1$, and $k_2$ be the incident, reflected, and transmitted wave vectors, respectively. The reflected amplitude can be written as

$$r = E_1/E_0 = -\frac{n-1}{n+1}. \quad (1.55)$$

If we write the reflectivity as

$$r = |r|e^{i\phi}, \quad (1.56)$$

then

$$\frac{r_+}{r_-} = \frac{r_+}{r_-}e^{i(\phi_+ - \phi_-)}. \quad (1.57)$$

The above relation says that the linearly polarized light will be converted to elliptically polarized light upon reflection because the circular components (RCP and LCP) will not have equal amplitude. The major axis of ellipse will
be rotated from the original direction of polarization of the incident light because of the phase introduced between two circular vibrations. The Kerr rotation and ellipticity can be expressed as

$$\theta_K = \frac{1}{2} (\phi_+ - \phi_-) \quad (1.58)$$

and

$$\varepsilon_K = \frac{r_+ - r_-}{r_+ + r_-}. \quad (1.59)$$

Now expressing the refractive index as a complex quantity, the complex Kerr rotation can be written as

$$\theta_K = \frac{2}{1 - \varepsilon} \left[ -(n_- - k_+) + i(n_+ - n_+) \right]. \quad (1.60)$$

The real component of $\theta_K$ gives the Kerr rotation and imaginary component of $\theta_K$ gives the ellipticity. The Kerr rotation depends upon the circular dichroism, whereas Kerr ellipticity depends upon circular birefringence of the medium. It can also be shown qualitatively that the Kerr rotation is proportional to the magnetization of a ferromagnetic material.

### 2.6.4 MOKE Geometries

There are three kinds of MOKE geometries, namely the longitudinal, transverse, and polar geometries depending upon the orientation of magnetization vector with respect to the sample surface and plane of incidence of light [90]. In the polar geometry the magnetization ($M$) lies perpendicular to the sample surface and parallel to the plane of incidence (Fig. 1.6a) whereas in longitudinal geometry $M$ lies parallel to the sample surface and the plane of incidence (Fig. 1.6b). The Kerr rotation can be explained

---

**Figure 1.6** The schematic diagrams of (a) polar, (b) longitudinal, and (c) transverse MOKE are shown.
qualitatively by considering the interaction of the electric field of the light with the magnetization of the material [91]. In a linearly polarized light (let say \( p \)-polarized), the electrons in the sample oscillate along the \( E \)-field of light. For a \( p \)-polarized light the oscillation is in the plane of incidence of the beam and also in plane of the sample. Again, a regularly reflected light experiences a \( \pi \)-phase change with respect to the incident beam. Therefore, the direction of \( E \) becomes opposite to the direction of the incident electric field. The Lorentz force on the oscillating electrons \((R_K = -m \times E)\) (Fig. 1.6a) generates an additional small vibrational component perpendicular to the plane of incidence. The electric field of reflected light is the vector sum of original electric field vector and Lorentz field. Figure 1.6a shows that the vector sum of the reflected lights, that is, the resultant electric field of reflected light, is rotated due to the magnetization of the sample. \( S \)-polarized light (electric field is perpendicular to the plane of incidence) gives similar Kerr rotation for the polar effect. The longitudinal Kerr effect is not observed at normal incidence as the cross-product is zero. For other incident angles the longitudinal Kerr effect is observed for both \( p \)- and \( s \)-polarized lights (Fig. 1.6b). The Kerr rotation \((\theta_K)\) and ellipticity \((\varepsilon_K)\) are related to each other by a relation: \( \theta_K + i\varepsilon_K = k/r \), in the limit \( k \ll r \) [92].

The third MOKE geometry is called the transverse geometry where \( M \) lies in the plane of the sample, but perpendicular to the plane of the incidence of light (Fig. 1.6c). Only \( p \)-polarized light shows transverse Kerr effect. In this case, the reflected beam remains linearly polarized without any Kerr rotation, but the amplitude is changed as magnetization vector changes sign from \(+M\) to \(-M\).

3. BACKGROUND OF TRMOKE

3.1. Background of TRMOKE Experiments

Advances in optics and invention of pico- and femtosecond lasers have brought about a drastic change in the measurement of ultrafast dynamics of matter. This new trend also got transferred into the spin and magnetization dynamics in the magnetic materials, particularly to cater the need to understand the fast dynamics in magnetic storage and memory units as well as for the understanding of the fundamental magnetic properties at the extreme time scale. Performing FMR experiments in the time domain and to verify the LLG equation to extract the precession frequency, damping and to directly image the time evolution of magnetization were the primary motivations for developing such a technique. In 1960s, time-resolved
magnetization dynamics was studied for the first time by digital storage oscilloscope and magnetic field pulse-induced magnetization oscillation with time period in the nanosecond time scale was observed [93–95]. In 1991, Freeman et al. [96] reported the first picosecond TRMOKE measurement of magnetization dynamics of magnetic thin film. The experimental technique was based upon magneto-optic sampling. A parallel strip coplanar transmission line structure is fabricated on a semi-insulating semiconductor substrate. This device is mounted on a chip carrier and is electrically connected to an external bias source. An above bandgap optical pulse is incident at the end of the biased transmission line creating transient photoconductivity in the semiconductor and launches a fast current pulse through the transmission lines [97]. Due to the opposite signs of the current pulses on the two sides of the lines a homogeneous magnetic field pulse is created in the gap between the lines. This yields a bandwidth of the order of 1 THz. The transient magneto-optic response of a sample overlaid or placed under the transmission lines can be measured at the region between the transmission lines by another pulse (probe) time delayed with respect to the pump pulse. By scanning the time delay between the pump and probe pulses by an optical delay line, the time-dependent magnetic response of the sample is measured with a time resolution similar to the pulse width of the laser. The rise-time of the pulse is determined by the carrier mobility and the optical pulse width, while the fall-time is governed by the carrier lifetime due to recombination and sweep-out in the presence of the bias field. All of these factors may be tuned over a broad range of values by the choice and preparation of the semiconductor, and by the excitation conditions. Using this technique, magnetization evolution and relaxation dynamics in pure and Tb-doped EuS thin films were measured at temperature down to $T=2.1\text{K}$ [98]. The spin–lattice relaxation time becomes an order of magnitude shorter due to the spin–orbit coupling after Tb doping in the EuS film. In 1992, Freeman et al. [99] showed clear precessional dynamics in an YIG film subjected to an in-plane static bias field. The pulsed magnetic field within the gap between the transmission lines is perpendicular to the plane of the film and thereby exerts a torque on the spin systems, causing a free precession of the spin systems. The bias field-dependent resonant frequency was modeled by Kittel formula, yielding $g=2.00$ and $4\pi M_s=1730\text{G}$ for YIG at room temperature. Freeman et al. also studied the intrinsic dynamics of magnetic flux threading superconductors to elucidate the issue of pattern formation in the type I intermediate state. The sample consists of a 100 nm thick Pb film sandwiched between a lithographic field coil and a
thin-film EuS-based magneto-optic sensor. An eddy-current decay time of 4.25 ns is found, from which the normal-state conductivity is found to be $\sigma_n \approx 1 \times 10^8 \text{S. cm}^{-1}$, which is much higher than the measured dc conductivity. This is believed to be a manifestation of the inhomogeneous spatial distribution of scattering sites, which gives rise to a frequency-dependent conductivity.

An important milestone in the early development of TRMOKE measurement was the direct measurement of the conduction electron spin–lattice relaxation time ($T_1$) in gold film [100]. The demonstration of the technique to polycrystalline gold films, a system with weak magneto-optic response and fast relaxation, showed that the technique can measure times as short as a few ps and can be applied at room temperature, where alternative indirect measurement techniques would fail.

Under a simple free electron bands for the conduction electrons, upon application of the instantaneous pulse, the spin up and down subbands are shifted creating a nonequilibrium. The resulting nonequilibrium distribution in the vicinity of the Fermi energy $E_F$ relaxes by spin–flip scattering and the rate at which the equilibrium moment $M_0$ arises is $1/T_1$. Thus, measuring the time dependence of the onset of the Kerr rotation in response to the instantaneous application of a magnetic field is a direct measurement of the longitudinal spin relaxation time $T_1$. A single turn coil with mean radius 1.4 $\mu$m is used to produce a 0.1 T field with rise-time of $\sim$4 ps at its center where a gold dot with 1 $\mu$m radius and 260 nm thickness is patterned (Fig. 1.7).

The time-resolved magneto-optic Kerr measurement setup is also shown in the bottom panel. An in-plane dc magnetic field of 2.8 T was applied to the sample. A room temperature measurement of the time-dependent Kerr rotation for a gold sample is displayed in Fig. 1.7 and the spin relaxation time $T_1$ is determined by fitting the detected polar Kerr rotation, using the measured time-dependent field $H(t)$. For a paramagnetic material of magnetic susceptibility $\chi$ and longitudinal spin relaxation time $T_1$ subjected to a time-varying field $H(t)$, the component of magnetization along the field is [96]

$$\mathbf{M}(t) = \frac{\chi}{T_1} \int_0^t H(t')e^{t'/T_1} dt'.$$

(1.61)

A fit with Eq. (1.61) determines $T_1 = 45 \pm 5$ ps. Transmission electron spin resonance (ESR) measurements of gold foil samples give a discrepancy by two orders of magnitude between values of probability of spin scattering
Figure 1.7 Top panel: micrograph of microfabricated coil and sample. Bottom panel: schematic layout of the magneto-optic experimental arrangement. Abbreviations: PD, photodiode; P1, polarizer; P2, Thompson polarizing beam splitter; BS1, 10% beam splitter; BS2, 50% beam splitter; PR, polarization rotator; DL, delay line; MO, microscope objective; DA, differential amplifier; PC, photoconductive switch; S1, GaAs substrate; S2, SiO₂ substrate; V, bias voltage. Reprinted with permission from Ref. [100]. Copyright 1996 by the American Physical Society.
per scattering event for foil and thin-film samples, and this suggests that Mattiessen’s rule is not obeyed over broad variations of mean electron scattering time.

### 3.2. Laser-Induced Ultrafast Spin Dynamics

In 1996, Beaurepaire *et al.* [57] reported another novel work, in which they showed ultrafast demagnetization and the subsequent spin dynamics in a ferromagnetic Ni thin film induced by 60fs optical pulses. Two types of measurements were done. In the first type the full hysteresis loops in the longitudinal MOKE configuration was measured at various time delays between the pump and probe beams, while in the second type pump–probe transmission measurement was done without any magnetic field. A striking feature obtained from the first type of measurement is a decrease of the remanence (signal at zero field) $M_r$ when the pump is on. While from the second type of measurement for laser fluence of 7 mJ/cm$^2$ (Fig. 1.8) shows a rapid decrease of $M_r$ within 2 ps, followed by a relaxation to a long-lived plateau. This figure clearly shows that the magnetization of the film drops during the first picosecond, indicating a fast increase of the spin temperature. For negative delays $M_r$ does not completely recover its value measured in the absence of pump beam. On the other hand, the transient reflectivity of the Ni film presents a single contribution, which is extremum for $\Delta t = 260$ ps,

![Figure 1.8](image)

**Figure 1.8** Transient remanent longitudinal MOKE signal of a Ni(20 nm)/MgF$_2$(100 nm) film for 7 mJ/cm$^2$ pump fluence. The signal is normalized to the signal measured in the absence of pump beam. The line is a guide to the eye. *Reprinted with permission from Ref. [57]. Copyright 1996 by the American Physical Society.*
showing that the contribution of nonthermal populations is weak and that the thermalization time is $\tau = 260\,\text{fs}$. From the differential transmittance measurements, the extracted electron temperature shows a sharp peak up to $T_e = 675\,\text{K}$ and then an exponential-like decay with a characteristic time of the order of $1\,\text{ps}$ until a saturation at $T_e = 550\,\text{K}$ is reached. On the other hand, the spin temperature rises rapidly during the first picosecond, presents a broad maximum at $T_s = 580\,\text{K}$ around $2\,\text{ps}$, and then, for longer delays, follows the same trend as the electron temperature. The authors presented a phenomenological model describing the heat transfer between three coupled baths (electrons, spins, and lattice). From the modeling the electron–lattice coupling constant $G_{el}$ is found as $8 \times 10^{17} \,\text{Wm}^{-3}\,\text{K}^{-1}$, electron–spin coupling constant $G_{es}$ as $6 \times 10^{17} \,\text{Wm}^{-3}\,\text{K}^{-1}$, and spin–lattice coupling constant $G_{sl}$ as $0.3 \times 10^{17} \,\text{Wm}^{-3}\,\text{K}^{-1}$. This work initiated the fundamental understanding of the dynamics of electron–spin interactions and would have important consequences on fast data writing technology. This was followed by the demonstration of full demagnetization in CoPt$_3$ film of 48.4 nm thickness by the same group in 1998 [101]. This effect looks similar to the ferro- to paramagnetic phase transition observed in static measurement, occurs in a time scale ($\sim 100\,\text{fs}$) when the electrons and spins are not in equilibrium with the lattice. Therefore, most of the demagnetization occurs during the thermalization of electron populations to a Fermi–Dirac distribution. The author suggested that the ultrafast demagnetization can be explained, taking into account the interactions between the spin-polarized electron population following the optical absorption. However, development of a microscopic theory of the spin dynamics during the time scale where the relaxation of electronic population occurs remained an open problem. In 2000, Koopmans et al. [102] raised questions over the observed magneto-optics in Ni—whether the signal is purely magnetic or has optical contributions. The authors claimed to have made a complete identification by explicitly measuring the time-resolved Kerr ellipticity and rotation, as well as its temperature and magnetic field dependence in epitaxially grown (111)- and (001)-oriented Cu/Ni/Cu wedges with Ni thickness ranging from 0 to 15 nm. In the first hundreds of femtoseconds the response is dominated by state filling effects. The true demagnetization takes approximately 0.5–1 ps. In the sub-ns time scales the spins are found to precess in their anisotropy field. Despite this debate experimental and theoretical works on ultrafast demagnetization continued to grow and formed an important subfield of magnetization dynamics. In 2002, van Kampen et al. [103] showed a novel all-optical method of excitation and detection of spin waves in 7 nm thick polycrystalline Ni layer on silicon with a special canted
configuration as shown in Fig. 1.9a. When the pump pulse heats the material at \( \Delta t = 0 \), a sharp decrease in \( M_z \) is observed. This effect is caused by a change in magnitude of the (temperature-dependent) magnetization. The subsequent recovery of \( M_z \) on a time scale of a few ps is due to rapid heat diffusion into the substrate. Strikingly, long after returning to thermal equilibrium a secondary response appears as a persistent oscillation that lasts for hundreds of picoseconds (Fig. 1.9b). The canting angle \( \theta_c \) is determined by the subtle balance between the external field and the anisotropy field of the film, including the shape anisotropy. Upon sudden heating by the pump pulse, not only the magnetization decreases but also the anisotropy of the film changes (Fig. 1.9c). This results in a change of the equilibrium orientation from \( \theta_c \) to \( \theta'_c \) (IIa), triggering an initial precession of the magnetization around its new equilibrium orientation (IIb). Heat diffusion into the film quickly removes the excess heat, and, after about 10 ps for metallic films, the original equilibrium angle is...
restored. However, at this point the magnetization is still not in equilibrium due to its initial displacement, and will continue to precess for hundreds of picoseconds (III). The progress in laser-induced magnetization dynamics was described in details in a review by Koopmans [104]. The research on laser-induced magnetization dynamics, in particular the ultrafast magnetization and different relaxation mechanisms, continued over the years [66,67,75,105–117], but are beyond the scope of this review.

In 1996 Guarisco et al. [118] reported time-resolved magneto-optical measurements of thermomagnetic writing in Gd$_{13}$Tb$_{12}$Fe$_{75}$ amorphous film on glass substrate and Co/Pt multilayer film on Si substrate. Curie-point thermomagnetic writing has been performed in less than 2 ns with the Co/Pt multilayer, whereas with the amorphous film magnetization reversal could not be observed within 6 ns. These differences in switching times are thought to be the consequence of the different thermal conductivities of the substrates. In 1997, Hiebert et al. [119] reported the first TRSKM study of the nonuniform precessional dynamics in a permalloy microelement described in details in Section 4. In the same year Stotz and Freeman [120] reported a high-resolution scanning optical microscope based upon solid immersion lens (SIL) for stroboscopic time-resolved studies of magnetic materials. They used both hemispherical (H-SIL) and the truncated-sphere solid immersion lenses (T-SIL). Calculation showed that in both cases, the same maximum $NA$ that can be achieved as $NA_{\text{max}} = n$, where $n$ is the refractive index of the material of the lens. Subsequently, $NA$ of up to 1.87 and spot size down to 0.22 $\mu$m were achieved. A thin-film magnetic recording head is used as a test specimen with well-defined optically reflective and magnetic features on a sub-$\mu$m length scale that allows characterization of the lateral spatial resolution of the instrument. In addition, the versatility of the microscope is demonstrated through examples of time-resolved magneto-optic imaging of the head. Further enhancement of the spatial resolution through the incorporation of confocal imaging is also described.

### 3.3. Imaging of Propagating Spin Waves

From the end of 1990s the TRMOKE experiment has become more popular to the magnetics community and a number of groups have come up with experimental demonstration of time-resolved magneto-optical measurements of precessional dynamics in ferromagnetic thin films and microstructures [121–127]. However, in 2003 Fassbender [128] introduced a new method to measure the time-domain propagation of magnetostatic spin
waves in iron garnet thin films. This method allows the determination of the phase of the spin wave as opposed to its complementary techniques like FMR and BLS. This is done by exciting spin waves by a short field pulse and detecting the propagating spin waves by time- and space-resolved magneto-optic Kerr magnetometry. The iron garnet film is 8 mm long, 2 mm wide, and 1.5 μm thick to show spin-wave propagation effects and is positioned face up on top of a 50 μm wide microstrip transmission line with the short axis parallel to the transmission line. A static magnetic field of 40 Oe is applied along the film width and the propagating spin waves are excited by magnetic field pulses of 1 Oe magnitude with 100 ps rise-time, 3 ns duration, and 200 ps fall-time along its length. Figure 1.10 shows the magneto-optic images of the central sample area for different delay times (in nanoseconds) after magnetic field-pulse excitation. At 0.4 ns after the start of the magnetic field pulse, the magnetization located directly above the microstrip transmission line is deflected most, as indicated by the white strip. This deflection leads to the formation of two spin-wave packets on each side of the transmission line, which propagate away from the localized field region to both sides. By evaluating the positions of bright and dark lines in successive images, which are chosen to correspond to half a precession period, the group velocity is determined as $v_{gr} = 6.6 \pm 0.4$ cm/μs. The arrows in Fig. 1.10 indicate the positions of the center of gravity of the propagating spin-wave packets and correspond to the group velocity.

To determine the phase velocity, the magneto-optic signal averaged across the width is plotted as a function of the distance from the transmission line along its length for different delay times, which gives the temporal evolution of constant phases. From their slopes, the phase velocity is determined as $v_{ph} = 34 \pm 7$ cm/μs.

3.4. Time-Resolved Kerr Effect Measurement on Samples Deposited on Opaque Substrate

An important development of the magnetic field pulse-induced TRMOKE experiment was brought about in 2002, when Barman et al. [129] showed delivering pulse to a spin valve fabricated on opaque substrate by the use of a transmission line fabricated on a 0.17 mm thin glass substrate and the dynamic Kerr rotation was measured through the transparent substrate of the transmission line as described in details in Section 5.1. A further development of this technique was demonstrated by Keatley et al. [130] where they used a hybrid coplanar stripline structure containing indium tin oxide windows in Au to perform optical FMR measurements on a sample grown
on an opaque substrate, using a pulsed magnetic field of any desired orientation. The phase of the oscillatory Kerr response was observed to vary as the probe spot was scanned across the coplanar stripline structure, confirming that the orientation of the pulsed field varied from parallel to perpendicular relative to the plane of the sample. A schematic of the typical experimental setup used in the above studies along with the transmission line and photoconductive switch arrangement and the configuration of the pulsed magnetic field with respect to the samples are shown in Fig. 1.11.
In 2006, Barman et al. [53] used an all-optical time-resolved Kerr effect measurement based upon a two-color collinear pump–probe technique to measure the picosecond precessional dynamics of single nanomagnets. The nanomagnets were optically pumped by linearly polarized strong laser pulses (\(\lambda = 400\,\text{nm}\), pulse width 100 fs). This causes an ultrafast demagnetization, following which an internal anisotropy field pulse is created in the system, which results in a change in the equilibrium magnetization orientation and triggers a precession. Linearly polarized weak laser pulses of 800 nm wavelength were time delayed with respect to the pump beam to probe the magnetization dynamics by detecting the polar magneto-optical Kerr rotation. The

![Figure 1.11](image_url)

**Figure 1.11** (a) A schematic of the time-resolved magneto-optical Kerr microscope, where the sample is excited by a magnetic field pulse. (b) The photoconductive switch and transmission line arrangement for generation of an optically triggered fast magnetic field pulse. Either the sample is overlaid onto the transmission line or a transmission line fabricated on a transparent substrate is overlaid on the sample to deliver the magnetic field pulse to the sample. (c) A schematic of the focused probe laser spot on a sample experiencing a pulsed magnetic field delivered by the transmission lines. The reflected beam path towards the detector is also shown.

### 3.5. All-Optical TRMOKE Microscope

In 2006, Barman et al. [53] used an all-optical time-resolved Kerr effect measurement based upon a two-color collinear pump–probe technique to measure the picosecond precessional dynamics of single nanomagnets. The nanomagnets were optically pumped by linearly polarized strong laser pulses (\(\lambda = 400\,\text{nm}\), pulse width 100 fs). This causes an ultrafast demagnetization, following which an internal anisotropy field pulse is created in the system, which results in a change in the equilibrium magnetization orientation and triggers a precession. Linearly polarized weak laser pulses of 800 nm wavelength were time delayed with respect to the pump beam to probe the magnetization dynamics by detecting the polar magneto-optical Kerr rotation. The
time-resolved magnetization dynamics was acquired point by point by varying the optical delay between the pump and probe and by measuring the Kerr rotation at each time delay. The pump beam was chopped at 2kHz and a phase-sensitive detection using a lock-in amplifier was used. Pump and probe beams were focused down to sub-μm diameter spots and precisely overlapped at the center of the sample at normal incidence using a single microscope objective. The back-reflected pump and probe beams are collected by a 50:50 beam splitter and are sent toward the detector. Before falling on the detector a small part of both pump and probe beams are steered into a CCD camera for viewing. A while light illumination is also sent to the sample through the same microscope objective for viewing the samples. Before the detector a spectral filter is kept to eliminate the pump beam from entering into the detector. Later, Pal et al. [131] and Rana et al. [132] have reported all-optical TRMOKE setups with improved detection system consisting of an optical bridge detector (balanced photodiode detection) and better spatial resolution and sensitivity. The optical bridge detector consists of a polarized beam splitter (PBS) and two photodiodes. The difference in the signal between the two photodiodes is proportional to the Kerr rotation [133]. The PBS is placed at 45° to the reflected light so that when a linearly polarized light (in absence of Kerr rotation) passes through the PBS, the intensity of light in two orthogonal components of polarization is identical that gives rise to a “balance” in the bridge. The Kerr rotation modifies the intensities in the two orthogonal components of polarization and gives rise to a finite electronic signal at the output of the optical bridge detector. For some samples such as Ni, Kerr ellipticity is much larger than Kerr rotation and in these cases the ellipticity is converted into rotation by introducing a λ/4 plate before the analyzer. A schematic of the typical all-optical TRMOKE microscope used by these authors is shown in Fig. 1.12. A typical configuration of the pump and probe beams focused through the same microscope objective and incident on the sample (an array of nanodots shown in Fig. 1.12) is also shown. The authors also used a λ/4 plate in front of the detector in some cases to measure the Kerr ellipticity instead of the Kerr rotation.

3.6. Benchtop TRMOKE Magnetometer

In 2008, Barman et al. [90] reported the development and application of a compact benchtop TRMOKE experiment based upon a picosecond pulsed laser. The excitation is done by a picosecond pulse generator, while the detection is done by an electronically synchronized picosecond laser pulse.
The system is integrated in a conventional upright microscope configuration with separate illumination, imaging, and magneto-optical probe paths. The system offers high stability, relative ease of alignment, sample changing, and a long range of time delay up to 40\,ns. Figure 1.13 shows a complete schematic of the benchtop TRMOKE magnetometer used for measurements of precessional dynamics of a permalloy microwire and a permalloy microdisk with a vortex state. The gyration mode of a vortex core was measured in the time domain over a time window of about 30\,ns. Later, the same technique was used to measure the gyration mode splitting of a vortex core [134] due to the magnetostatic interactions from other disks in the array as described in details in Section 8.

4. TIME-RESOLVED IMAGING OF MAGNETIZATION DYNAMICS IN MICROSCOPIC MAGNETIC ELEMENTS

With the increasing speed of operation of the magnetic data storage and other emerging technologies, fundamental interest in microscopic origin of the magnetization dynamics in confined magnetic elements has
become increasingly more important. TRMOKE has been considered as the natural candidate for investigation of such phenomena due to the intrinsic simultaneous spatiotemporal resolution associated with any optical technique.

4.1. Imaging Nonuniform Precessional Dynamics in Single Ferromagnetic Microstructures

In 1997 Hiebert et al. [119] reported the first direct observation of time-domain evolution of nonuniform magnetization precession in a lithographically patterned circular permalloy element with 8μm diameter by TRSKM. The dynamics was excited by pulsed magnetic field created by a biased GaAs-photoconductive switch sending a current pulse around a lithographic gold coil. The transient magnetic field is created at the center of the coil, perpendicular to the plane of the substrate (Fig. 1.14). A time-delayed probe beam is used to probe and image the dynamics with a spatial resolution better than 0.7μm. Measurement of the time-resolved Kerr rotation with the laser spot placed at the center of the element showed well-developed damped precession of magnetization. The bias

![Figure 1.13](image_url) (a) The schematic of the complete benchtop TRMOKE magnetometer. The dotted arrow shows proposed arrangements, while the solid arrows correspond to the existing arrangements. (b) The synchronization scheme between the pump and probe pulse is shown. Reprinted with permission from Ref. [90]. Copyright 2008 by the American Institute of Physics.
Figure 1.14 Images showing (a) the spatial dependence of the Kerr rotation for different times at $H_\parallel = 500$ Oe, along with schematic illustration of two possible mechanisms underlying the nonuniformity; (b) demagnetizing effects at the edges carrying free poles; and (c) asymmetry of the magnetic vector phase trajectory. Reprinted with permission from Ref. [119]. Copyright 1997 by the American Physical Society.
field-dependent precessional frequency is fitted with the Kittel formula, while the full time-domain data were modeled by numerically solving the LLG equation and by considering the profile of the pulsed magnetic field obtained by measuring the time-resolved Kerr rotation at a very high bias magnetic field. At such a large bias field the precession amplitude becomes negligibly small and the magnetization follows the envelope of the pulsed magnetic field. The temporal shape of the pulsed field is found to be somewhat irregular due to reflections generated by the transmission line and its contact with the transmission line. The damping coefficient $\alpha$ is found to be 0.008, which is close to the value obtained by other method [135]. In order to account for a visible divergence of the oscillation amplitude for intermediate values of time between the experimental and modeling data, the authors have acquired scanning Kerr images of magnetization at selected time delays at a bias field $H = 500$ Oe. The scanning images clearly showed spatially nonuniform dynamics with the initial evolution of Kerr rotation near the edges carrying free magnetic poles, which eventually propagate toward the center. A number of possible reasons for the nonuniform excitation of dynamics in the element are discussed but reasons like eddy current and thermal effects have been ruled out and the primary reason for the observed nonuniformity is found to be the nonuniform demagnetizing fields near the edges of the elements carrying free magnetic poles. This concept is strengthened by the observed symmetry in the response, which is repeated with the sample orientation rotated by 90º, which also confirmed the negligible role of the in-plane magnetic anisotropy. In addition, possible nonuniformity in the pulsed magnetic field and nonlinear process were also considered to have some effects on the observed nonuniform dynamics. Later the same group [136,137] used the same technique to also image the magnetization reversal, and characterizations of the speed of magnetic recording devices.

In 2000, Acremann et al. [123] from ETH Zurich reported the first three-dimensional imaging of precessional orbits of the magnetization vector in a magnetic field using a time-resolved vectorial Kerr experiment that measures all three components of the magnetization vector with picosecond resolution. The time-resolved dynamics of a flat, polycrystalline Co disk of 6 $\mu$m in diameter, grown inside a single turn aluminum coil was measured. The equilibrium magnetic configuration measured by spin-polarized scanning electron microscopy showed a flux closure domain pattern, suggesting negligible magnetocrystalline and other sources of magnetic anisotropy for this disk. Measurement of two components ($M_y$ and $M_z$) of time-resolved
Kerr rotation from a location where the initial spin configuration is along \( x \)-direction shows that in agreement with Larmor’s theorem, \( M_y \) deviates from the equilibrium value before \( M_z \) does, and remains ahead of \( M_z \) by approximately \( \pi/2 \). Spatially resolved magnetization vector at a fixed time delay (180 ps) shows that a field pulse along \( z \)-direction activates the \(+y (-y)\) component of the magnetization within domains with equilibrium spin configuration along \(+x (-x)\) and leads to the magnetic contrast observed in the in-plane components. However, despite being symmetric with respect to rotations around the \( z \)-axis, \( M_z(t) \) reveals a spatial nonuniformity observed also by Hiebert et al. Apparently the nonuniform excitation is a wave propagating from the border toward the center and then bouncing back but the observed propagation velocity is an order of magnitude lower than the spin waves with the expected \( k \) vector and an apparent sign reversal as opposed to what is expected for a Neumann boundary condition, where the wave should be reflected without changing sign. Further calculation of the spatial dependence of the magnetic field produced by the coil reveals a nonuniformity that can be as large as 20% across the disk. On the other hand, the observed time offset indicates that the origin of the nonuniform precessional mode is some delay in the start of the precession possibly due to the Faraday’s induction law. The calculated delay of 10 ps corresponds to an estimated rise-time of the magnetic field pulse of about 10 ps. Although small corrections due to eddy current is important, after the initial delay of 10 ps the eddy current is no longer important due to the large precessional period of \(~150\) ps and the experimental geometry.

### 4.2. Imaging of Noise by Time-Resolved Kerr Microscopy

So far it was believed that the stroboscopic imaging of magnetization dynamics is limited to repetitive phenomena as the signals are averaged over trains of pulses instead of coming from a single event. However, in 2000 Freeman et al. [138] reported a method for extracting information corresponding to underlying stochastic behavior through spectroscopic analysis of random telegraphic noise in the signal. The only required condition for the extraction is that the measurement system noise level should be lower than the level of random noise associated with stochastic processes in the sample. Development of low noise and stable solid-state lasers enabled a shorter integration times by more than an order of magnitude. Consequently, stochastic changes in the sample are manifested in the signal with much greater statistical significance. The technique was applied to a
magnetic recording head and the noise associated with the random switching behavior was imaged by time-resolved Kerr microscope. The interpretation of the observed behavior is not straightforward because an essentially two-dimensional imaging method (only with vertical depth of ~ optical penetration depth) is used to image a three-dimensional device. However, the observed noise structure is reproducible from run to run on the same head, but varies dramatically from device to device, and is absent for devices reversing in a completely repetitive manner within the experimental resolution. Later, results from micromagnetic modeling of thermally assisted switching in a small platelet are presented to help elucidate the procedure and its interpretation [139].

4.3. Imaging of Spin-Wave Modes in Ferromagnetic Microwires

In 2002, Park et al. [140] reported the direct observation of localized spin-wave modes in individual thin-film permalloy wires with 2 and 5 \( \mu \)m widths using TRSKM. Full spatiotemporal imaging of the dynamics were performed in the BWVMS and DE geometries by compiling successive snapshots of the cross section of the microwire at fixed time steps of 20 or 40 ps. In the DE geometry the edge modes (EMs) do not appear. In the BWVMS geometry frequency domain images show evolution of both the center and edge modes [141]. The confinement is strongest at the highest magnetic fields. In the case of the 2 \( \mu \)m wire at 400 Oe, the edge and center mode (CM) images overlap in space, although the edge mode appears as a weak satellite of the dominant center mode peak. As the magnetic field decreases, the frequencies of both the center and edge modes decrease and the spatial extent of the edge modes grows while that of the center mode shrinks. In the case of the 2 \( \mu \)m wire, the modes merge below 75 Oe into a single mode spanning the entire wire. The same process occurs in the 5 \( \mu \)m wire at a lower magnetic field. At the lowest fields, the single mode observed in the 2 \( \mu \)m wire actually increases in frequency as the external field decreases. The allowed modes were numerically reproduced by employing WKB argument of Schlömann and Joseph [142]. This method approximates the equation of motion for the magnetization vector by the same form as a Schrödinger equation, and bound states can then be found from the Bohr–Sommerfeld quantization condition:

\[
\int_{X_1}^{X_2(v)} q(v,H(x)) \, dx = n\pi.
\]  

(1.62)
The boundary for the outer turning point $x_1$ is determined as the position where the sum of the applied and demagnetization fields vanish, and is computed by treating the poles at the edges of the wire as magnetic charges. The second turning point $x_2$ is the position corresponding to the internal field at which there is no longer a real wave vector associated with the frequency. Practically, it is found by searching for the field that has a minimum in the dispersion relation at a given frequency, and then using the calculated field profile to find the turning point $x_2$. Using similar method the same group has imaged the spin dynamics of closure domain and magnetic vortex structures [143], which is described in greater details in Section 8. In 2004, Bayer et al. [144] imaged the magnetization dynamics of inhomogeneously magnetized stripes with 2 $\mu$m width and 18 nm thickness. When the in-plane bias magnetic field applied parallel to the short axis of the stripe is on the order of the shape anisotropy field, the equilibrium magnetization near the stripe edges rotates $90^\circ$ over a length scale of order $100 \text{nm} - 1 \mu$m. Using a combination of semianalytical theory and micromagnetic simulations, it was found that these modes span the entire stripe but can only be detected near the edges, where the effective wave vector is small. This experiment demonstrates that the inhomogeneous system has well-defined collective modes, including higher frequency excitations that are accounted for in the spin-wave model.

In 2008, Kruglyak et al. [145] studied low-amplitude precessional dynamics within and near the demagnetized region at one end of a 200 $\mu$m long and 6 $\mu$m wide Co$_{80}$Fe$_{20}$(4 nm)/Ni$_{88}$Fe$_{12}$(10.8 nm) wire. The variation in the frequency of the observed modes is governed by the local value of the internal static effective magnetic field and the local angle between the static magnetization and the effective direction of the wave vector of the confined modes. Dynamic images showed irregular stripes formed by the dynamic magnetization perpendicular to the long axis of the wire. This transient stripe pattern is interpreted in terms of a collective mode of the quasiperiodic system of ripple domains, excited by a non-uniform torque due to the uniform pulsed field acting upon nonuniform static magnetization.

### 4.4. Configurational Anisotropy in Precession Frequency and Damping

In 2003, Barman et al. [146] reported imaging of magnetization precession in a square permalloy element with $10 \times 10 \mu$m$^2$ area and with 150 nm thickness. The dynamics was excited by a pulsed magnetic field, while a bias
magnetic field ($\mathbf{H}$) is applied at different directions within the plane of the element. The precession frequency showed a strong fourfold anisotropy superposed on a weak uniaxial anisotropy with the variation in the bias field orientation angle. Modeling of the precession frequency gives a fourfold anisotropy field (Eq. 1.16) $4K_4/M$ as $-33\text{Oe}$, while the uniaxial anisotropy (Eq. 1.15) is about $2\text{Oe}$. The authors also observed an anisotropy in the apparent damping behavior where the damping occurs much faster when $\mathbf{H}$ is applied parallel to a diagonal of the element as opposed to that when $\mathbf{H}$ is applied parallel to one of the edges of the element. The pulsed field profile was measured by electro-optic sampling by using a LiNbO$_3$ crystals placed on top of the transmission line structure and by measuring the electro-optic Kerr rotation with the laser spot placed at the center of a track and at the center between two tracks [147]. The rise-time of the pulse is measured as about $35\text{ps}$, while some secondary peaks due to the reflections of the current pulse in the transmission line are observed. The stronger damping observed for $\mathbf{H}$ parallel to the diagonal is not the result of this coherent pumping mechanism due to the presence of these secondary peaks in the pulsed field. This is clear since the enhanced damping is observed for a range of different field values as the phase of the precession is varied relative to that of the secondary field peaks. Later, measurement of damping as a function of the bias field orientation was performed, which showed a complete fourfold variation of damping with the bias magnetic field [148]. The time-resolved Kerr images were acquired for $\mathbf{H}$ applied parallel to the edge (Fig. 1.15a) and diagonal (Fig. 1.15c) of the square to understand more details of the apparent damping. The dynamic images have shown that this is associated with spatial nonuniformity at the center of the element. When $\mathbf{H}$ is applied parallel to the edge of the element, the nonuniformity is confined to the edges until longer delay times. This work demonstrates that the shape of a thin-film element can have a major influence upon the apparent damping of the magnetization precession.

In a further work, the authors characterized the eigenmode in greater details [149]. For an element of 10$\mu$m width, and with an optical resolution of better than 1$\mu$m, one may observe spin waves with wave vector $k$ in the range $6 \times 10^3$–$6 \times 10^4\text{cm}^{-1}$. The frequencies of these magnetostatic waves depend upon the product of $k$ and the film thickness. For a continuous film of 150$\text{nm}$ thickness and a static field of about 1kOe, the frequencies of the BWVMS waves lie more than 1GHz below that of the uniform mode. The frequencies of the excited modes were determined by measuring the transient Kerr rotation with the probe spot placed 1, 2, 3, and 4$\mu$m away from
the center of the element on the symmetry axis parallel to $\mathbf{H}$. As the probe spot is moved away from the center, two new lower frequency modes appear whose intensities increase away from the center. These new modes appear consistently for a number of bias magnetic fields. For $\mathbf{H}$ applied parallel to the diagonal of the element five modes are observed as the probe spot is placed at the center of the element and moved away from the element along a straight line parallel to the direction of $\mathbf{H}$. With $\mathbf{H}$ applied parallel to an edge, only the highest frequency mode was found to have significant amplitude at the center of the element, whereas a number of modes had significant amplitude at the center when $\mathbf{H}$ was applied parallel to a diagonal. Dephasing hence leads to a larger apparent damping at the center of the element when $\mathbf{H}$ is applied parallel to a diagonal as compared to when $\mathbf{H}$ is applied parallel to an edge. Later, the dynamic images of that element for two different bias field geometries were reproduced by micromagnetic simulations (Fig. 1.15b and d). Subsequently, the spatial character of the resonant modes was simulated by applying an ac field to the sample.

Figure 1.15 (a, c) Experimental and (b, d) simulated time-resolved magneto-optical Kerr effect images at various delay times after excitation of a square-shaped Ni$_{81}$Fe$_{19}$ microelement by a picosecond pulsed magnetic field. The static spin configurations (column 6), internal magnetic field map (column 7), and the geometries of the applied bias field are shown at the right end of the figure. The difference in the evolution and the subsequent dephasing of the spin waves for the two different geometries leads to an anisotropic damping in this study. Adapted from Ref. [150]. Copyright 2007 by the American Institute of Physics.
oscillating at a particular mode frequency [150]. This revealed that for $H$ parallel to the edge of the element, the highest frequency mode has significant amplitude over the entire area of the sample, apart from the narrow demagnetized regions near the edges. However, the lower frequency modes were found to be localized away from the center of the element and closer to the edge of the sample as their frequency decreased. For $H$ applied parallel to the diagonal of the element, all five modes were found to be localized. The lowest frequency mode was most strongly localized and situated closest to the center of the element. The spatial width of the modes increased with frequency and the positions of peak amplitude shift away from the diagonal toward the corners. This work clearly demonstrated that dephasing of the resonant mode spectrum must be considered as a primary extrinsic source of damping of precessional motion in microscale elements.

This work was further extended to investigate the dependence of fourfold magnetic anisotropy and the anisotropy of the apparent damping upon shape and aspect ratio in micron-sized permalloy elements [151]. The circular element showed no fourfold anisotropy or anisotropic damping, while all the square elements showed fourfold anisotropy and in some cases also an anisotropic damping. The fourfold anisotropy field and the anisotropy in the apparent damping increase with decreasing aspect ratio. The anisotropy in the apparent damping increases with the number of fringes observed within the element. The number of fringes is related to the time taken for the excited mode spectrum to dephase which is in turn determined by the aspect ratio of the element.

4.5. Excitation and Imaging Individual Resonant Modes by Time-Resolved Kerr Microscopy

A new development in the TRSKM was attempted in 2008 by Neudert et al. [152] by introducing phase-locked harmonic excitation of individual resonant modes instead of a broadband excitation by a pulsed magnetic field. Using this technique they identified and imaged two resonant modes in a 40 micron wide and 160 nm thick $\text{Fe}_{70}\text{Co}_{8}\text{B}_{12}\text{Si}_{10}$ square element. However, excitation with the frequencies of the individual modes resulted in a superposition of both modes although with different amplitude. By detecting the out-of-plane and in-plane Kerr rotations, they investigated the relative shape of the magnetization orbit. The orbit does not change its relative shape at different positions relative to the edge of the element.
4.6. Imaging Large Angle Reorientation of $2 \times 2 \mu m^2$ CoFe/NiFe Element

Keatley et al. [153] studied a large angle reorientation of a $2 \times 2 \mu m^2$ CoFe/NiFe element with 4.6 nm thickness by applying a large amplitude (59 Oe) pulsed magnetic field with 150 ps rise-time generated by a pulse generator. Analysis of the experimental data revealed that the magnetization reoriented by $10^\circ \pm 10^\circ$ but only at the center of the element, while the canted magnetization near the edge regions remained pinned. A large observed damping (0.1) was suggested due to the excitation of spin waves in the regions of nonuniform magnetization.

4.7. Imaging Magnetization Dynamics of Hard Disk Writers

In 2011, Gangmei et al. [154] have shown direct application of time- and vector-resolved Kerr microscopy at the wafer level on magnetic hard disk writers. Pulsed field excited magnetization dynamics within the yoke and pole piece of a partially built writer with a three turn coil and a single layer CoFe yoke of 200 nm was studied by measuring the in-plane (parallel and perpendicular to the bias field) and polar Kerr rotations. The response times of different magnetization components at different locations has been related to the orientation of the static magnetization, the torque generated by the coil windings, and the propagation of flux through the yoke and pole piece. The Kerr images showed localization of the perpendicular magnetization component, along the short axis of the yoke, and its propagation into the confluence region where it presumably magnetizes the bridge region. These images provide direct experimental confirmation of the well-discussed “flux beaming” mechanism for the operation of the yoke. The flux beaming occurs along the short axis because the magnetization at the edges of the yoke is restricted to rotate in response to a pulsed field for the minimization of demagnetizing energy. More recently Yu et al. [155] have reported magnetization dynamics of a multilayered writer yoke using a similar measurement scheme as Ref. [154]. Dynamic images showed that while the yoke takes stripe domains in absence of bias field, the equilibrium state of the tip during pulse cycling is metastable and depends on the polarity of the driving current. Consequently different combinations of bias field and current polarity lead to different dynamics deep within the confluence region, which is also suggested by recent modeling work [156].
4.8. Conversion of Free Space Microwave to Magnonic Architecture: TRSKM Imaging

In 2011, Au et al. [157] demonstrated a method of coupling uniform FMR-based antenna to global free space microwave and converting them directly to finite wavelength spin waves propagating in stripe magnonic waveguide. The idea is based upon the difference in the frequency of the uniform precessional mode in the stripe and the continuous film and the broken translational symmetry at the junction between the continuous film and the end of the stripe—the later allowed the uniform mode in the film to overcome the momentum gap, as required for the excitation of finite wave vector spin waves in the stripe. Subsequently, the authors imaged the propagating spin waves in the stripes using time-resolved Kerr microscope and modeled the results to extract the wavelength of the spin waves as 36 \(\text{m} \mu\text{m}\). Later in 2012, the same authors demonstrated a new magnonic source based upon resonant enhancement of microwave field near a resonating magnetic nanowire transducer [158]. In their scheme, a number of transducers located at different points of the magnonic chip are driven into uniform precession by a global uniform microwave field. Each of the transducers generates a rotating stray magnetic field due to the dynamic magnetic charges induced on its lateral surface due to the uniform precession. The stray field is strongly localized and this localization enables the coupling of the stray field to spin waves of short wavelength. The multiple magnonic sources on the chip are driven by the same microwave clock signal and therefore they maintain a constant phase relationship.

A number of other applications of time-resolved scanning Kerr microscope have been reported such as imaging precessional switching (Section 10), magnetic vortex and domain wall dynamics (Section 8), and coherent suppression of small amplitude precession (Section 9), which are described in dedicated sections in this review.

5. TIME-RESOLVED MAGNETIZATION DYNAMICS OF MAGNETIC MULTILAYERS

Over the last few decades, there has been a rapidly increasing interest in magnetic multilayers due to their novel properties like exchange bias [159,160], perpendicular magnetic anisotropy (PMA) [161], giant magnetoresistance (GMR) and spin valve [162], tunnel magnetoresistance (TMR) [163], exchange spring magnets [164], and spin–torque effects [165,166]. In addition to the fundamental science involved they have existing or
potential applications in magnetic storage, memory, and other spintronic devices. For faster operation of the devices it is important to understand how these properties behave at the ultrafast time scales. However, the initial study of picosecond magnetization dynamics of magnetic multilayers was not straightforward due to technical issues of delivering a fast excitation field pulse to the sample fabricated on Si or other opaque substrate.

5.1. Time-Resolved Magnetization Dynamics of Spin Valves and Exchange-Coupled Bilayers

The first report on time-domain study of precessional dynamics in magnetic multilayer system appeared in 2002 when Barman et al. [129] employed a new technique of delivering excitation field pulse to the spin valve deposited on an opaque substrate by a transmission line made on transparent and thin glass substrate, which can be overlaid on the sample. A biased-photoconductive switch was connected to this transmission line, which was optically triggered to create the excitation pulse, while the dynamics was probed in between the transmission lines through the glass substrate. A series of Si(100)/Ta(50 Å)/Ni_{81}Fe_{19}(50 Å)/Cu(d Å)/Co(50 Å)/IrMn(100 Å)/Ta(30 Å) where \( d = 10 \text{ Å} \) (S1), 20 Å (S2), and 30 Å (S3) for three samples were studied. While S1 showed a single precession frequency, S2 and S3 showed two modes for all bias fields applied along the easy and hard axis of the spin valves. The variation of precession frequency with the bias field and the magnetic hysteresis loops were simultaneously modeled by solving LLG equation for a trilayer system with an interlayer coupling term in the magnetic free energy. The modeling gives magnetic parameters including the anisotropy and magnetization of Co and Ni_{81}Fe_{19} layers, magnitude and direction of the exchange bias, and the interlayer coupling. The interlayer coupling is found to be an order of magnitude higher in S1 as compared to S2 and S3. In the presence of interlayer coupling, the two modes correspond to collective excitations of the trilayer, in which the layer magnetizations precess either in or out of phase. The magneto-optical response from the two layers may cancel for the out-of-phase mode, particularly in the presence of strong ferromagnetic coupling as in S1. This technique was meant for using in wide variety of samples, in particular to magnetic read–write devices at their intermediate steps of fabrication.

An all–optical pump–probe experiment on the temperature dynamics in exchange-biased bilayers was reported by Weber et al. [167]. Excitation by a 8.5 ps laser pulse created a hot spin and phonon gas in a 5 nm thick Ni_{81}Fe_{19} (F) and a 10 nm thick Fe_{50}Mn_{50} (AF) layer. The spin–lattice temperature
dynamics is sensed in real time by the time evolution of the exchange bias field on the picosecond time scale. A calibration with temperature-dependent quasistatic Kerr measurements yields a pump–pulse–induced temperature increase of about 100°C at the interface. Upon photoexcitation, the exchange coupling across the interface between the ferromagnetic and antiferromagnetic layer is reduced within the first 10 ps, leading to a reduction of the bias field to about 50% of its initial value. The fast thermal unpinning of the exchange coupling is followed by a heat-diffusion-dominated recovery with a relaxation time on the order of 160 ps. The thermal diffusion length and the effective diffusivity have been evaluated for the whole layer system. The measured time evolution of the exchange bias anisotropy can be interpreted as an effective internal anisotropy pulsed field, which can induce coherent precessional magnetization dynamics in the ferromagnetic layer.

Engebretson et al. [168] studied the reversal mechanisms in exchange-biased bilayers FeF₂/Fe. While many systems showed some form of reversal asymmetry, in the transition metal difluoride antiferromagnets the reversal occurs largely by rotation on the left side of the hysteresis loop and by nucleation of reverse domains on the right side. They studied the time-resolved dynamics of the out-of-plane component of the magnetization induced by an ultrafast magnetic field pulse with temporal width of about 120 ps and with magnitude of about 5 Oe. The sample was placed with the ferromagnetic film on top of a tapered microstripline mounted on the cold finger of a helium flow cryostat to study the temperature dependence. An increase in the characteristic precession frequency with decreasing temperature was observed, even above the Néel temperature of the antiferromagnet, which was interpreted in terms of the anisotropy enhancement due to antiferromagnetic spin fluctuations. Below the Néel point the magnetization precession is strongly suppressed due to the damping provided by exchange coupling to the antiferromagnetic layer. Dynamic hysteresis loops show a distinct asymmetry, which does not occur in the corresponding static loops. The asymmetry is due to the suppression of coherent rotation at the magnetic fields where reverse domain nucleation is known to be favorable. A nonmonotonic temperature dependence of the dynamical asymmetry is observed, which is not entirely understood but displays correlations with the temperature dependence of the coercivity.

5.2. Ultrafast Magnetization Dynamics of Magnetic Multilayers with PMA

Ever since the discovery of PMA, it has become an interesting system for bit-patterned media and spin transfer torque–magnetoresistive random access
memory (STT-MRAM) devices. The origin of PMA remained a subject of interest, and experimental and theoretical evidences suggest that strong and localized interfacial $d$–$d$ hybridization produces an enhanced perpendicular Co orbital moment responsible for the PMA. In the early 2000, significant efforts have been made in understanding their magnetization reversal mechanisms. The reversal is strongly dependent on the number of bilayer periods $n$. At the thick film limit, multilayers with $n > 10$ nucleation of reverse domains from isolated defects occur, which then develop into a dendritic domain pattern; while at the other extreme Pt/Co/Pt trilayers showed magnetic domains rapidly expanding by propagation of smooth and well-defined domain walls. Intermediate $n$ values show an unusual reversal behavior with domain formation on two different length scales. However, the ultrafast dynamics of the PMA multilayers has become increasingly important to address the write speed issues in such systems. In addition the damping coefficient $\alpha$ is also of importance to have a lower write current in the STT-MRAM devices, while retaining the high thermal stability.

In 2007, Barman et al. [169] reported the first experimental study of time-resolved magnetization dynamics in $[\text{Co}(4\,\text{Å})/\text{Pt}(8\,\text{Å})]_n$ multilayers with $n = 2, 5, 8$, and 12. The static Kerr magnetometry provided square hysteresis loops typical of the high PMA samples with a significant variation in the shape of the loops with $n$. All samples showed an ultrafast demagnetization within 400fs and biexponential decay with 10ps and 400ps decay constants, irrespective of $n$. A fast precession with frequencies up to 81 GHz was observed, which was identified as the uniform precession mode of the whole multilayer stack. A significant decrease in precession frequency is observed with increasing $n$. The extracted PMA increases with reduced $n$ from 12 down to 5 and becomes almost constant below $n = 5$. The reduction of PMA with increased $n$ may be associated with an increase in interface roughness or a variation in the crystal structure of the Co layer. The damping coefficient $\alpha$ ($\geq 0.1$) increases systematically with the increase in $n$ but remained independent of the bias field. The authors attributed the large damping coefficient to the enhanced spin–orbit coupling in ultrathin magnetic films, while its enhancement with $n$ was thought to be due to an enhanced interaction between magnons and conduction electrons.

Mizukami et al. [170] reported the magnetization dynamics of Pt/Co ($d_{\text{Co}}$/Pt films with $0.5 \leq d_{\text{Co}} \leq 2.0$, where static polar MOKE measurements showed PMA for multilayers with $d_{\text{Co}} \leq 1.0$. The precession frequency $f$ and relaxation time $\tau$ were measured as a function of angle $\theta_H$.
between field and direction normal to the film, varied from 0° to 80°. The experimental angular-dependence data of \( f \) and \( 1/\tau \) were fitted with expressions derived from the LLG equation \( f = f_0 \sqrt{1 - (2\pi f_0 \tau)^{-2}} \), where \( f_0 = (\gamma/2\pi)\sqrt{H_1 H_2} / \sqrt{1 + x^2} \) and \( 1/\tau = x\gamma(H_1 + H_2)/(1 + x^2) \). Here, \( H_1 = H \cos(\theta_H) + H_{k\text{eff}} \cos 2\theta \) and \( H_2 = H \cos(\theta_H - \theta) + H_{k\text{eff}} \cos 2\theta \). The fitting treated \( \gamma \) and \( x \) as fitting parameters, while using \( H_{k\text{eff}} \) as obtained from the static polar MOKE (PMOKE) experiment. The extracted \( K_{u\text{eff}} \) is found to be linearly proportional to \( 1/d_{Co} \) in accordance with the static measurements, whereas \( x \) is not proportional to \( 1/d_{Co} \), which is also confirmed from FMR measurements. The authors also extracted the relaxation frequency \( G = x\gamma M_s \), which is linearly proportional to \( 1/d_{Co} \) for \( d_{Co} > 1.0 \text{nm} \) with a slope of \( 34 \times 10^8 \text{rad/s/mm} \), which is three times larger than the Pt/Py/Pt film. However, analysis showed that this enhancement is not due to the spin pumping effect but possibly due to the Co 3d–Pt 5d hybridization effect, which decreases the bandwidth \( W \) for the Co atomic layer in contact with a Pt layer, enhancing both PMA and Gilbert damping. However, this hybridization could not explain the significant increase in \( G \) for \( d_{Co} < 1.0 \text{nm} \). This thickness regime is close to the interface roughness or alloying layer thickness, which might affect Gilbert damping but the problem remained open. Later in 2011, the same authors reported [171] systematic investigations of \( x \) for Pt-buffered and –capped Ni/Co multilayer films [Si/SiO2/Ta(5)/Pt(2)/Co/[Ni/Co]_n/Pt(2), units in nm] with different Co (Ni) layer thicknesses \( d_{Co} \) (\( d_{Ni} \)) and bilayer period \( n \) while maintaining the ratio \( d_{Ni} / d_{Co} = 2 \). The \( x \) and \( K_{u\text{eff}} \) depended strongly on the layer thickness and bilayer periodicity and rise to maximum values of about 0.08Merg/cm³ and 8Merg/cm³, respectively. Damping rates exhibited linear dependence with respect to inverse multilayer thickness, indicating that \( x \) in Ni/Co multilayers is enhanced mainly at interfaces in contact with Pt layers and does not increase much inside multilayers. In the same year Pal et al. [131] reported magnetization dynamics in [Co(\( t_{Co} \))/Pd(0.9nm)_8] MLs with \( t_{Co} \) varying between 1.0 and 0.22nm. A combination of static PMOKE and vibrating sample magnetometer (VSM) measurements showed that all samples show PMA, while the anisotropy field increases systematically with decrease in \( t_{Co} \) and exhibits a maximum at 0.22nm, beyond which it decreases sharply. The saturation magnetization decreases monotonically with the decrease in \( t_{Co} \) over the entire range down to 0.13nm. X-ray reflectivity measurements showed an average roughness at the interface to be about 0.05nm. The time–resolved Kerr rotation showed a single precession frequency, which varied by a large amount
from 5 to 92 GHz with the decrease in $t_{Co}$. The extracted PMA from macrospin modeling showed a linear variation with $1/t_{Co}$ (Fig. 1.16). The damping coefficient $\alpha$ was extracted from the relaxation time and precession frequency $f$ as $\tau = 1/2\pi f\alpha$ and was also found to be linearly proportional to $1/t_{Co}$ (Fig. 1.16). Consequently, $\alpha$ was found to be linearly proportional to $K_{\text{eff}}$ with a slope of $4.33 \times 10^{-8}$ cc/erg. Discussion on possible mechanisms on the enhancement of $\alpha$ ruled out magnon–magnon scattering as for perpendicularly magnetized samples it is less effective. The spin pumping effect is also found to be very small from the slope of the plot of relaxation frequency

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**Figure 1.16**  (a) The bias field dependence of the experimental precession frequencies (symbols) and the calculated frequencies (solid line) with Kittel formula are shown for multilayers with different Co layer thickness. (b) The geometry for the macrospin model is shown. (c) The damping coefficient $\alpha$ (symbols: experimental data, solid line: linear fit) is plotted as a function of $1/t_{Co}$. (d) The extracted perpendicular magnetic anisotropy $K_{\text{eff}}$ and the saturation magnetization $M_s$ (filled squares: from TRMOKE, open circles: from static magnetometry) are plotted as a function of $t_{Co}$. The dashed line shows the calculated $M_s$ values, while the dotted line corresponds to the linear fit to $K_{\text{eff}}$ versus $t_{Co}$. Reprinted with permission from Ref. [131]. Copyright 2011 by the American Institute of Physics.
G with $1/t_{Co}$. The third possibility could be interface roughness and alloying, which would increase $\alpha$ but would also decrease $K_{eff}$, which is opposite to their observation and hence this possibility is also ruled out.

Subsequently, they found that $d$–$d$ hybridization at the Co/Pd interface leads to the simultaneous increase in $\alpha$ and PMA. Liu et al. [172] also reported ultrafast magnetization dynamics of $[\text{Co}(t)/\text{Pd}(9\,\text{Å})]_8$ multilayers with $t = 2.8$, $5.0$ and $7.5\,\text{nm}$. They also observed a dramatic increase in the magnetization precession frequency and about 2.6-fold increase in damping constant with the reduction in $t$. They used a simplified micromagnetic model optimized for magnetic multilayers to qualitatively and quantitatively reproduce the experimental data. However, dependence of $\alpha$ on PMA was not discussed possibly due to the lack of sizeable amount of data points.

Pal et al. [173] reported dipole-exchange spin-wave spectra in a series of CoO-capped $[\text{Co}(t)/\text{Pt}(7\,\text{Å})]_{n-1}\text{Co}(t)$ multilayer systems, where the total Co moment $(n \times t)$ is constant. The Co layer thickness within the ML is varied from $0.2\,\text{nm}$ ($n = 40$) to $0.8\,\text{nm}$ ($n = 10$) in this experiment. Since the experiments are performed at room temperature the CoO layer can be considered paramagnetic with no exchange bias effect. However, the CoO top layer still introduces an asymmetry between the top and bottom of the Co/Pt ML. The experimental spectra in general showed two intense peaks and additional lower intensity peaks, which can no longer be modeled by the macrospin model. The spin-wave spectra were calculated in a linear approximation, by numerically solving the Landau–Lifshitz equation on the discrete lattice—discrete dipole approximation (DDA) method [174], which is limited to calculating the standing wave modes. The calculation showed that the spin-wave spectra consist of low-frequency surface waves (antisymmetric and symmetric), and a band of the bulk spin waves. For a symmetric surface, the two surface waves are degenerate, but the presence of two different materials Pt in the bottom and CoO on the top resulted in an asymmetric surface, which splits the degeneracy between the symmetric and antisymmetric surface modes in their case.

### 6. PRECESSIONAL DYNAMICS OF MAGNETIC NANODOT ARRAYS

Experimental study of magnetization dynamics in magnetic nanodots and its arrays has started to appear in the literature in 2002. The initial works were done primarily by FMR [175–177] and BLS spectroscopy [178,179]. The FMR measurements showed dipole-exchange spin-wave dispersion for
a perpendicularly magnetized sample where in-plane wave vector is quantized due to the finite dot radius. The inhomogeneity of the intradot static demagnetization field in the nonellipsoidal dot found to be important for the dynamics. The BLS measurements primarily showed standing spin-wave modes of Damon–Eshbach origin, dipole-exchange mode and reminiscent of backward volume mode, in addition to a laterally confined edge mode, with its frequency independent of dot radius. Time-domain magnetoresistive method was used to study the spin-transfer-induced magnetization reversal in current-perpendicular spin-valve nanomagnetic junctions by using a pulsed current bias [180]. Later, spin-torque-driven magnetization precession modes and damping have been studied and electrical control of damping of nanomagnet was proven by Krivorotov et al. [54] by magnetoresistive method.

6.1. Size-Dependent Crossover to Nonuniform Precession

Time-resolved magneto-optical study of magnetization dynamics in nanodots was reported not earlier than 2005, when Kruglyak et al. [181,182] reported the pulsed field-induced precessional dynamics in arrays of magnetic nanodots with varying diameter from 630 nm down to 64 nm. The Co$_{80}$Fe$_{20}$$(10\text{Å})$/Ni$_{88}$Fe$_{12}$$(27\text{Å})$ dots were prepared by EBL and ion milling, while the interdot separations were not systematic and only the dependence of the size of the dot on the magnetization dynamics was studied here. The frequency spectra of the time-resolved Kerr rotation shows that the precession frequency increases with the decrease in element size, while the same trend continue as the element size reduces from 120 to 64 nm, although with an overall reduction of frequency. In the 220 nm element array, the higher frequency modes continue the trend observed for the 630 and 425 nm element arrays, while the low-frequency modes continue the trend of the 120 and 64 nm element arrays. This means a crossover from one branch of modes to another at element size close to 220 nm (Fig. 1.17). Generally several branches are present in the spectra but their relative amplitudes are different and at 220 nm two branches have equal amplitudes.

Micromagnetic simulations showed that the observed two branches of modes correspond to the higher frequency center mode and the lower frequency edge mode. The demagnetized regions near the edges result in lower effective magnetic field and lower frequency for the edge mode. The crossover occurs due to the increased relative area of demagnetized regions and consequently the relative contribution of the edge mode with the reduction
in the element size. Simultaneously the relative localization area of the center mode decreases and its contribution also decreases. The most striking feature is that the dynamics becomes more nonuniform with the decrease in element size, which may have significant implications in the noise of magnetic read sensor operating by precessional switching. Further numerical extension of this work [183] showed the appearance of additional modes in the spin-wave spectrum as the thickness of the element is increased. The frequency of the edge mode is sensitive to the strength of the exchange interaction, dipolar interactions with nearest neighbor elements, and rounding of the corners of the elements and edge roughness. Extensive experimental and simulation studies were performed by Keatley et al. [184] with variable element size (separation) of 637(25), 428(17), 236(77), 124(30), and 70(37) nm but with larger thickness of Co$_{80}$Fe$_{20}$(40 Å)/Ni$_{88}$Fe$_{12}$(108 Å). This higher thickness leads to greater nonuniformity in static magnetization and the total effective field within the element. Experimentally, they observed the coexistence of two branches of excited modes above a particular bias field. Below the crossover field, the higher frequency branch disappears. Micromagnetic simulations showed that the higher frequency branch have high mode amplitude at the center of the element in regions of positive effective

**Figure 1.17** Extracted mode frequencies are plotted against element size for different bias fields. The symbols represent data points and the curves are simulations. A small horizontal shift has been applied to the experimental frequencies at different bias fields for clarity. The error bars represent the full width at half maximum of the spectral mode peaks. *Reprinted with permission from Ref. [181]. Copyright 2005 by the American Physical Society.*
field while the lower frequency branch have high mode amplitude near the edges of the element perpendicular to the bias field. The crossover between the higher and lower frequency branches was mediated by the complicated evolution of the total effective field within the element. Below the crossover region edge type mode extends into the entire element. Simulations revealed that the majority of the modes are delocalized with finite fast Fourier transform (FFT) magnitude throughout the element and this delocalized nature of the excited modes seems to be an intrinsic property of nanoscale nonellipsoidal elements. Collective nature of the modes with various spatial distributions over the arrays was also observed.

6.2. Dynamical Configurational Anisotropy

Kruglyak et al. [185] studied the dynamical configurational anisotropy in $4 \times 4 \mu m^2$ arrays of square elements of $220 \text{nm}$ width and $2.5 \text{nm}$ thickness and with $95 \text{nm}$ edge-to-edge separation by TRMOKE excited by a pulsed magnetic field. The size of the elements ensures that they do not possess static configurational anisotropy [186]. Both the number and frequency of the precessional modes show fourfold anisotropy with the variation of the bias field orientation, where two modes are observed for field parallel to the edge and a single mode for field parallel to the diagonal of the element. Analysis with micromagnetic simulation shows that the observed anisotropy is governed by the variation of both the static and dynamic magnetization configuration and the associated dynamic effective magnetic field and is therefore termed as dynamic configurational anisotropy. More detailed analysis showed that the dynamic configurational anisotropy can be qualitatively explained in terms of the angular dispersion of dipolar-exchange spin waves. The changes in the effective wave vector with the bias field orientation drive the variation in frequency of the lowest lying mode, although a rigorous quantitative analysis is required to understand the complete nature of the eigenmodes.

6.3. Coupling Between Magnetic and Elastic Modes

A study of optically induced magnetization dynamics in patterned Ni nanodisks with $185 \text{nm}$ diameter and $280 \text{nm}$ separation was studied by Müller et al. [187]. The ultrafast demagnetization has no effect on nanostructuring, while the transient reflectivity is enhanced due to the plasmonic resonance during pump-pulse duration. On the other hand the GHz frequency spin-wave dynamics is significantly altered in the...
nanostructures. The thickness dependence of the excited magnetic mode allowed to quantitatively study the dipole interaction between single Ni nanodisks. In addition, coupling between magnetic and elastic modes could not be proven or ruled out from this study.

6.4. Collective Spin-Wave Dynamics in Arrays of Nanomagnets

Precessional dynamics is found to be drastically different in a coupled nanodot array as compared with an isolated single nanodot. The dynamics of an array cannot be interpreted in terms of frequencies of individual dots rather a new set of normal (collective) modes appear. This is very important to understand as the limit of the areal density of the bit-patterned media and magnonic devices.

6.4.1 Imaging Collective Magnonic Modes

In 2010, Kruglyak et al. [188] reported the measurement and imaging of collective magnonic modes in arrays of $80 \times 40 \text{nm}^2 \text{Co}_{50}\text{Fe}_{50}(0.7\text{nm})/\text{Ni}_{92}\text{Fe}_{8}(4.5\text{nm})$ elements with $\sim20\text{nm}$ edge-to-edge separation by time-resolved Kerr microscopy. The FFT spectra of the time-resolved Kerr rotation excited by an in-plane pulsed magnetic field measured at different values of in-plane bias magnetic field shows a broad single peak at higher bias field and splitting into three narrower peaks at smaller bias field values. Micromagnetic simulations on arrays of $3 \times 3$ elements reproduced the collective nonuniform precessional modes such as quasiuniform, backward volume-like and Damon–Eshbach-like modes (Fig. 1.18). Since the modes are detectable, they must have a wavelength greater than the optical spot size, which was demonstrated by imaging the collective modes by exciting the array with harmonic fields with each particular collective mode frequency. The experimental images present complex spatial pattern as opposed to the clear collective modes observed in the simulated images, possibly due to the presence of reminiscent of the neighboring modes in each collective mode. There are possibility of the existence of a number of closely spaced non-uniform collective modes, which were not resolved in the experiment. In conclusion, such arrays behave like a quasicontinuous “magnonic” metamaterial, with properties defined by the geometry of the array that supports long wavelength collective spin waves.

6.4.2 Effect of Areal Density on Collective Magnonic Modes

In 2011, Rana et al. [132] used an all-optical TRMOKE microscope to excite and detect collective magnetization dynamics in arrays of $200\text{nm}$
wide square permalloy dots with 20 nm thickness. They used arrays with systematically varying areal density where the edge-to-edge separation ($S$) changes from 50 to 400 nm. Although the ultrafast demagnetization and remagnetization do not vary with the areal density the precessional dynamics vary significantly. A single precessional mode observed for $S = 50$ and 75 nm splits into three modes at $S = 100$ nm and with the increasing value of $S$ the splitting amplitude increases. At $S = 400$ nm only two modes are observed. At the same time when the bias field is varied for the sample with $S = 50$ nm, the single mode splits into two or three modes as the bias field is reduced to 0.85 kOe and below. Detailed micromagnetic simulations [189] showed that a single dot of the studied dimensions will have a CM, a mixed edge-Damon–Eshbach (EM-DE) mode and an EM. In the array with $S = 50$ nm the profile of the single mode observed in the experiment and simulation are neither of the three modes observed in the single element but a uniform collective mode extended over almost the entire area of the dots. The calculated stray magnetic field in between the dots is >2 kOe or more, which makes the whole array to behave almost like a continuous film, where the elements lost their individual identities. Similar trend is retained for the array with $S = 75$ nm. For $S = 100$ nm, the highest frequency mode (mode 1) corresponds to the collective precession of the elements where the mode profiles of individual elements look similar to the CM of a single element. In mode 3, the power profile of individual element looks identical to the mixed EM-DE

**Figure 1.18** Simulated images of the magnitude (top) and phase (bottom) for the modes at 4.5, 4.8, and 5.4 GHz of the $3 \times 3$ array are shown for a bias field of 197 Oe. Reprinted with permission from Ref. [188]. Copyright 2010 by the American Physical Society.
mode of single 200nm element, whereas mode 2 is BWVMS like mode. This trend continues for $S$ up to 300nm but for $S = 400$nm, the mode profiles of individual elements look similar to the modes observed for a single 200nm element. For this separation, the dots are magnetostatically isolated and collective behavior is not observed. Extraction of damping shows that for $S = 50$nm, the effective damping coefficient is about 0.017, similar to that measured from a thin film of same thickness, but increases to about 0.025 at $S = 100$nm and almost saturates beyond that up to $S = 400$nm. This again shows a transition from a strongly collective to weakly collective dynamics in these dot arrays. The authors investigated an anisotropy in the collective dynamics with the variation of the azimuthal angle of the bias magnetic field [190]. The collective dynamics is found to be anisotropic with the variation in the azimuthal angle of in-plane bias field. As the azimuthal angle of the bias field is changed from 0° to 45° with respect to the symmetry axis a systematic transition from a uniform collective dynamics to non-collective dynamics is observed in strongly collective regime ($S = 50$nm), whereas no clear trend is observed in the weakly collective regime ($S = 100$nm). The frequencies of the modes also showed a fourfold anisotropy, which were modeled by using Eq. (1.16) and the fourfold anisotropy fields obtained from experimental and simulated data agree well. As the size of a single dot is reduced, the lowest frequency mode (EM) becomes more dominant and its frequency increases, while the higher frequency modes (CM and EM-DE) loose power and finally the 50nm dot shows a single peak (EM).

The same authors have reported an all-optical time-domain excitation and detection of ultrafast magnetization dynamics of arrays of lithographically patterned 50nm permalloy dots down to the single nanodot regime [191]. A collinear pump–probe arrangement was used for this measurement, which enabled to achieve a very good spatial resolution and sensitivity even in an all-optical excitation and detection scheme of the precessional dynamics. In the single nanodot regime the dynamics reveals one dominant resonant mode of the magnetization corresponding to the precessional dynamics of magnetization of the edge regions of the nanodots. The precession of magnetization of the central region of the nanomagnet was found to be extremely weak as opposed to the popular perception. The resonant mode has slightly higher damping (0.023) than that of the unpatterned thin film (0.017). With the increase in the areal density of the array the precessional frequency increases due to the increase in the magnetostatic interactions (Fig. 1.19).
A simple modeling shows that quadrupolar interaction is dominant over dipolar interaction for these square elements. The damping also increases with the areal density due to the dynamic dephasing of precession of a number of weakly interacting dots (Fig. 1.19). The apparent damping showed a sudden jump at $S = 50\text{nm}$ due to the superposition of two closely spaced modes. Numerical simulations show that the higher frequency mode is a coherent precession of the EM in the array, while the new lower frequency mode is a collective BWVMS like mode of the entire array.

![Graph](image)

**Figure 1.19** (a) The precession frequency is plotted as a function of width to center-to-center distance ratio ($w/a$). The circular and square symbols correspond to the experimental and simulated results, respectively, while the solid curves correspond to the fit to $f = f_0 - A(w/a)^3 + B(w/a)^5$, where $A$ and $B$ are the strengths of the dipolar and quadrupolar interactions. (b) The damping coefficient $\alpha$ is plotted as a function of $S$. The symbols correspond to the experimental data, while the solid line corresponds to a linear fit. The dashed line corresponds to the measured value of $\alpha$ for a continuous permalloy film grown under identical conditions. *Reprinted with permission from Ref. [191]. Copyright 2011 by the American Chemical Society.*
6.4.3 Effect of Lattice Symmetry on Collective Magnonic Modes

In 2013, Saha et al. [192] have reported the fabrication and time-domain measurements of magnetization dynamics in two-dimensional arrays of closely packed circular nanodots arranged in different lattice symmetry, namely square, rectangular, hexagonal, honeycomb, and octagonal symmetry. The dots were chosen as circular to have no shape anisotropy, 100 nm in diameter to possess two precessional modes (EM and CM), and of 30 nm separation between nearest edges to have strong magnetostatic interaction and collective dynamics. The above conditions ensured that effects of inter-dot magnetostatic interaction arising purely from lattice symmetry will be studied. The optically induced precessional dynamics was measured by an all-optical TRMOKE microscope. The dynamics has also been simulated by a time-dependent micromagnetic simulation and the time-domain magnetization, power spectra, and the power and phase profiles of the resonant modes have been numerically calculated to analyze the dynamics. A significant variation in the mode spectra is observed with the variation of lattice symmetry (Fig. 1.20). The square lattice shows a single precessional mode, which is found out to be the coherent precession of the edge mode of all the dots in the lattice. The rectangular lattice shows two modes, which are

![Figure 1.20](image-url)  
Figure 1.20 (a) Scanning electron micrographs of permalloy nanodot lattices with varying lattice symmetry. (b) Time-resolved Kerr rotation. (c) Corresponding FFT spectra from different nanodot lattices. (d) Simulated spin-wave spectra for the permalloy nanodot lattices with $7 \times 7$ elements. Adapted with permission from Ref. [192]. Copyright 2013 by the Wiley.
coherent precession of the edge modes or the center modes of the dots over the entire lattice apart from the edges of the lattice. The hexagonal lattice shows three clear modes. They correspond to the uniform collective mode, a BWVMS-like mode and a bowtie-like mode of the entire lattice. The honeycomb lattice shows broad and rich spin-wave spectrum, which includes various localized and extended modes including the DE-like mode of the lattice. For octagonal lattice three modes are observed, which are uniform, DE-like, and localized modes of the lattice. In the hexagonal lattice, the edge mode of the dots forming a BWVMS-like collective mode of the lattice shows a clear sixfold anisotropy with the variation of the azimuthal angle of the bias field, while other modes do not show any clear anisotropy. Analysis of the power map of the anisotropic mode clearly shows a sixfold variation of the net dynamic component of magnetization with the azimuthal angle as a result of a similar variation of the magnetostatic field.

6.4.4 Effect of Shape

In addition to the physical structures of the individual elements, the magnetic properties of the nanomagnetic metamaterials in the form of artificial lattices may also be tailored by varying the shapes of the individual elements. The magnetization of confined magnetic elements deviates from the bias magnetic field even in presence of a large enough magnetic field thus creating demagnetized regions. These demagnetized regions play important roles in determining their magnetic ground states as well as the magnetization dynamics. They offer different potentials to the propagating spin waves and also spatially localize them within the demagnetized regions. In addition, the interdot magnetostatic interaction in ordered arrays of nanomagnets may also get affected by the shapes of the elements. This is because the profile of the stray magnetic field depends on the shapes of the boundaries of the elements as well as the internal magnetic field. The magnetization dynamics of nanomagnets have primarily been confined in square, circular, and ellipsoidal shapes. Configurational anisotropies originating from the shapes of the elements have been observed in square-shaped elements. However, time-resolved dynamics of nanomagnets of other shapes is rare. Recently, Mahato et al. [193] reported time-resolved magnetization dynamics in cross-shaped permalloy nanoelements of 600 nm length and width, 20 nm thickness, and about 50 nm separation between the nearest edges. The precessional dynamics is studied under the application of a bias magnetic field, whose orientation $\phi$ is varied with respect to the sample geometry. A significant anisotropy in the nature and frequencies of the
spin-wave modes with $\phi$ is observed. When the bias field is applied parallel to one of the arms of the cross (0° and 90°) a single resonant mode is observed. However, a band of modes appear as $\phi$ deviates from 0° to 90°. The band becomes broadest at $\phi = 45°$. The lowest frequency mode shows a fourfold anisotropy. Simulated mode profiles show the occurrence of a number of quantized modes, where the frequencies of the modes increase with the mode quantization numbers. The frequencies and profiles of the modes of the elements in the array are significantly affected by the neighboring elements as $\phi$ deviates from 0° and become maximum at 45°.

7. MAGNETIZATION DYNAMICS IN SINGLE NANOMAGNETS

The quest to measure and understand the intrinsic dynamics of single or magnetostatically isolated nanomagnet from its environment at various time scales has long been an interested problem, which intrigued the researchers. Such a problem demands development of techniques of simultaneous spatial and temporal resolution at the extreme level, which has been the primary hindrance in the development of this research field. Although X-ray microscopy [55] and time-domain magnetoresistance techniques [54,180] made some early progress, they have many drawbacks in terms of technical specifications as well as due to the requirements of special facilities and sample fabrication requirements. Being an optical technique TRMOKE has the intrinsic advantage in temporal resolution particularly due to the emergence of ultrashort laser pulses and an improvement in spatial resolution will certainly take the overall spatiotemporal resolution of this technique very high as compared to other alternative techniques.


The idea is based upon enhancing the magneto-optical Kerr rotation by coating a ferromagnetic material with a dielectric enhancement layer. By properly choosing the thickness and material index of the dielectric layer, the Kerr signal reflected off the magnetic surface can be enhanced through constructive multiple reflections. This idea was used to enhance static Kerr signal from unpatterned film [194–196] or from ferromagnetic nanodots [197]. However, application in the measurement of the picosecond magnetization dynamics from single nanomagnets came in 2006 by Barman et al. [53]. They measured cylindrical Ni dots with constant height of
150 nm and with varying diameter \((D)\) from 5 µm down to 125 nm. The dots were coated with 70 nm thick SiN layer to have a fivefold increase in the Kerr rotation due to cavity enhancement. The ultrafast magnetization dynamics in such nanomagnets was excited optically by a femtosecond laser pulse and was probed by an all-optical TRMOKE setup. The dynamics showed one or more resonant modes, where the highest frequency mode showed a nonmonotonic variation of frequency with the \(D\) (aspect ratio—AR). Between 5 µm (AR = 0.03) and 2 µm (AR = 0.075) the frequency remains almost constant. Below 1 µm (AR = 0.15) the frequency increases steadily with reducing diameter down to 300 nm (AR = 0.5). At \(D = 250\) nm (AR = 0.6) the frequency undergoes a sharp jump by 33%, which then increases at a steeper rate down to \(D = 125\) nm (AR = 1.2). This step is clearly a physical phenomenon as the frequency variation between different magnets of the same size is less than 1%. The authors modeled the variation of the precession frequency with dot diameter by solving LLG equation under small angle approximation by assuming shape, volume, and surface anisotropy energies, which explains the observed variation barring the sudden jump in the precession frequency at an AR = 0.6. A sudden change in the effective volume anisotropy constant from \(1 \times 10^4\) erg/cm\(^3\) for samples >250 nm to \(1.5 \times 10^5\) erg/cm\(^3\) for samples ≤250 nm was taken into account to explain this. Although the origin of such a drastic change is not entirely clear, evidence of such a change in Ni nanomagnets has been reported from quasistatic hysteresis loop measurements [198,199]. Later in 2007, the authors reported a size-dependent effective damping \((\zeta_{\text{eff}})\) of the same set of nanomagnets extracted directly from the time-resolved Kerr rotation data [200,201]. Large qualitative and quantitative differences in \(\zeta_{\text{eff}}\) between the micro- and nanoscale elements are observed (Fig. 1.21). For magnets with diameters between 5 and 3 µm, \(\zeta_{\text{eff}}\) remains almost constant at nearly 0.17. At 2 µm, \(\zeta_{\text{eff}}\) reduces sharply followed by a slow decrease down to 500 nm, where it attains a value ~0.04, comparable to the reported damping coefficient 0.05 of continuous Ni thin films measured by all-optical method [103]. For magnets with \(D > 2\) µm a strong bias field dependence of \(\zeta_{\text{eff}}\) is observed, while for magnets with \(D < 1\) µm, no bias field dependence is observed. For magnets with intermediate \(D\) weak bias field dependence is observed. Even for magnets with \(D > 2\) µm, \(\zeta_{\text{eff}}\) approaches smaller values at larger bias fields, nearly the value obtained for the nanoscale magnets. Further analysis showed that the dephasing of two closely spaced modes in the frequency domain opens additional dissipative channel and causes the enhancement of the effective damping of the uniform precession mode...
for samples with diameter $>1 \, \mu m$. For $D < 1 \, \mu m$ two frequency modes always remain well split for the whole bias field range down to 1.25 kOe, contributing no extrinsic sources to the damping of magnetization precession. The authors claimed that the observation of field-independent
damping for nanoscale magnets close to the thin-film value ensures the reliability of coherent control of precessional switching by a straightforward pulse shaping scheme in nanomagnets.

7.2. Ultrafast Thermal Switching, Relaxation, and Precession of Individual Ferromagnetic Disks

Laraoui et al. [73,74] reported ultrafast magnetization dynamics in CoPt$_3$ nanodisks and permalloy microstructures. The authors used a two-color collinear pump–probe technique using 790nm laser as pump and its second harmonic, 390nm as probe beams focused to 600 and 300 nm spots, respectively, for the measurement. They performed two types of measurements: first the imaging of the magnetization of the disk for a fixed pump–probe delay by scanning the sample mounted on a piezoelectric stage and second is the time-resolved magnetization dynamics and reflectivity in the center of the disk. The CoPt$_3$ disks fabricated on (0001)-oriented sapphire substrate are 15nm thick with diameter from 1µm down to 0.2µm, whereas square-shaped permalloy microstructures with 30µm width and 20nm thickness were fabricated on glass substrate. The dynamic magnetization images at a fixed time delay of 300fs was taken for the two adjacent disks. First the magnetizations of both disks were saturated to $M_s$ by applying a 4kOe field and then the field was reduced to $-100$ Oe. The scanned image of the first disk taken with a density of laser excitation of 1mJ/cm$^2$ shows a uniform dark contrast over the entire disk area. The second disk was first switched to $-M_s$ with a pump density of 8mJ/cm$^2$. The pump–probe image taken with a pump density of 1mJ/cm$^2$ showed a uniform bright contrast. This implies that for high enough density of excitation, the process of thermal switching induced by the pump pulse modifies the magnetization uniformly over the entire disk. The time-resolved data shows an initial demagnetization followed by a fast ($\tau_{\text{spin–lat}}=5.2$ ps) and a slow ($\tau_{\text{diff}}=530$ ps) remagnetization. The fast remagnetization is associated with the equilibrium between the spin and lattice, while the slow remagnetization corresponds to the thermal diffusion to the surrounding of the disk deposited on the sapphire substrate. The magnetization of the permalloy structure was measured with densities of excitation of 8 and 4mJ/cm$^2$ and by applying a 3.5kOe field perpendicular to the film surface. The time constant for remagnetization is found to be longer for larger densities of excitation ($\tau_{\text{spin–lat}}=11$ ps and 3.6ps, respectively) due to the temperature dependence of electronic specific heat as can be obtained from a 3TM. In both cases $\tau_{\text{diff}}=900$ ps, much higher than the CoPt$_3$ sample due to the less efficient energy transfer to the
glass substrate. In the case of permalloy a precessional dynamics is induced as opposed to the CoPt$_3$ sample and the period of precession increases significantly with the density of excitation. This is due to the density-dependent shape anisotropy, which reduces the effective field for higher density leading toward lower precession frequency.

### 7.3. Effects of Magnetic Ground States on Magnetization Dynamics in Single Nanodisk

An interesting feature of nanomagnets is its different magnetic ground states such as vortex and quasisingle-domain states under varying conditions depending on the minimization of total energy resulting from competing energy terms such as demagnetizing and exchange energies. Liu et al. [202] showed that the shape of the nanodisks significantly affects the vortex-to-quasisingle-domain state transitions. The authors studied pulsed field excited magnetization dynamics of isolated permalloy nanodisks of 160nm diameter by time-resolved Kerr microscope with a spot size of 700nm. The time-resolved out-of-plane component of magnetization showed different modal frequencies due to the transition between different domain states. At very high field the disk is in single-domain state ($H_0 \sim 100$ kA/m) with a dominant single frequency mode. As the field is reduced, it transforms into quasisingle-domain state and the variation of frequency with bias field follows a modified Kittel formula [203] including an effective demagnetizing factor $F(\beta)$ for a cylinder, where $\beta$ is the thickness/radius:

$$f_0 = \frac{g_0}{2\pi} \sqrt{H_0^2 + H_0 M_s \left[1 - 3F(\beta)\right]}.$$  \hspace{1cm} (1.63)

However, at low $H_0$, C-state configuration leads to significant lower frequencies than predicted by Eq. (1.63), which is based upon uniform spin configuration. As $H_0$ is lowered below 29kA/m, a frequency jump (increase) by 10GHz is observed due to a transition to a vortex state and the frequency remains almost constant at ~13GHz for $H_0$ is reduced to 0. As $H_0$ is increased beyond 29kA/m, frequency still remains constant till 40kA/m, beyond which frequency drops to ~11GHz. This lower frequency is characteristic of a thinner disk and suggests that the vortex core is affected by sloping disk edge. As $H_0$ is increased to 68kA/m, abrupt frequency drop by ~5GHz is observed registering a vortex annihilation and transition to a quasiuniform state. The measured vortex nucleation and annihilation fields corresponded within 1kA/m for repeated bias field sweeps.
Another interesting observation is that in a single cycle, transition can be spontaneously triggered when the bias field is kept fixed at a critical field region and the disk is driven into dynamics. This was demonstrated at two fixed bias fields 67.4 and 29.1 kA/m, which showed a vortex to quasisingle-domain transition and vice versa during the dynamics. The authors claimed that single nanodisk measurement is essential for these observations as measurement over an array would give a gradual transition over a small range of $H_0$ due to small physical variation between the disks in the array.

7.4. Large Amplitude Magnetization Dynamics in a Single Nanomagnet

Previous works showed that the small amplitude dynamics of a square nanomagnet magnetized in-plane and along an edge is complicated by excitation of center and edge modes [181,184,204]. The two modes coexist for a range of bias fields and dot size while maximum amplitude shifts from center to edge mode as the bias field or the dot size is decreased. However, large amplitude excitation can significantly modify the spatial character of the dynamics. Keatley et al. [205] studied pulsed field-induced large amplitude magnetization dynamics of a single $440 \times 440$ nm$^2$ CoFe(1 nm)/NiFe(5 nm)/CoFe(1 nm) nanomagnet by a TRSKM with a 300 nm spot size. The pulsed field has <40 ps rise-time, 70 ps duration, and up to 90 Oe of peak amplitude ($h_p$). Experimental spectra were obtained after applying bias field ($H_B$) applied parallel to the edge and diagonal of the element showed that a single mode is excited for all $H_B$ values. The spectra for $H_B$ parallel to the edges at two different pulsed field amplitudes (30 and 80 Oe) were modeled by micromagnetic simulations with reduced edge $M_s$ following a profile of the form, $\cos^6(2\pi x/l)\cos^6(2\pi y/l)$, where $x$ and $y$ are spatial coordinates and $l=440$ nm and with damping coefficient $\alpha=0.03$. The simulations showed that for $h_p=80$ Oe a remarkable suppression of the edge mode amplitude occurs for $H_B$ between 400 and 250 Oe. For $h_p=30$ Oe the crossover field shifts from 400 to 300 Oe, while for $h_p=80$ Oe, the crossover field shifts to 200 Oe, indicating that for larger values of $h_p$ the excitation of the center mode is favored at lower bias field values. The onset of edge mode suppression can be observed in the simulated spectra for an $h_p$ value as small as 10 Oe. The laser spot size (300 nm) is smaller than the permalloy square (440 nm) and hence highly localized edge modes at large $H_B > 500$ Oe are
probed by the Gaussian tail of the optical probe. Therefore, sensitivity to edge modes for these bias fields is diminished so that they were not observed experimentally. For $H_B$ smaller than 500 Oe but greater than the crossover field the edge mode extends more toward the central region of the element where it is more readily detected. However, weaker coupling of the edge mode to the large amplitude field and modification of the internal magnetic field lead to the suppression of the edge mode. Controlled suppression of edge mode has potential applications in nanoscale spin-transfer torque oscillators and bistable switching devices, for which amplitude of magnetization trajectory are generally large and more uniform dynamic response is desirable.

7.5. Detecting Single Nanomagnet Dynamics in Varying Magnetostatic Environment

Liu et al. [206] demonstrated a different approach of detecting single nanomagnet dynamics by creating an array where neighboring magnets have very different magnetization dynamics. Although the individual elements cannot be resolved optically using far-field detection, their dynamic response will be distinguishable from the FFT spectra. The studied 150nm thick Ni nanomagnets with 150nm diameter embedded in 500nm diameter magnets with edge-to-edge spacing between the 150 and 500nm magnets as 200nm (Fig. 1.22). The samples were fabricated on anti-reflection coated (for 800nm) Si substrates by EBL. They measured the time-resolved Kerr rotation from the individual 500nm, 150nm magnets and the array where the 150nm magnet is embedded by four 500nm magnets. They showed that the dynamics of the 150nm magnet can be identified from those of the 500nm magnets due to their large size difference, which gave large frequency difference ($\sim 15$ GHz at 4.5 kOe). The bias field dependence of the dynamics showed that while the frequency of the 500nm magnet does not vary much with the bias magnetic field, the 150nm magnet shows a large variation as was also obtained by micromagnetic simulations. To quantify the magnetostatic effects from the neighboring elements, a 150nm magnet in between two 500nm magnets with varying edge-to-edge separation from 200 to 100nm was studied. As the edge-to-edge separation decreases from 200 to 100nm, the edge mode frequency of the 150nm magnet decreases by about 5 GHz, due to the variation of magnetostatic field on that element from its neighbors. This shows that the effect of the neighboring magnetic elements cannot be ignored in the design and implementation of these systems.
8. MAGNETIC VORTEX AND DOMAIN WALL DYNAMICS

Submicron diameter magnetic disks form a flux-closure domain pattern with a core inside which magnetization rotates out of the plane direction. This is due to the competition between magnetostatic and exchange energies. Minimization of the magnetostatic energy by curling around the center comes at the cost of large exchange energy of the out-of-plane magnetized core.

Magnetization configurations of the vortex structure have been shown in Fig. 1.23 [5]. Inhomogeneously magnetized vortex state has two key
properties. The first is the polarity of the core \((p)\) with \(p = \pm 1\), and the second is the chirality \(c = \pm 1\). Chirality represents the sense of rotation, which is clockwise or anticlockwise whereas polarity refers to the direction of the out-of-plane core magnetization (up or down). The dynamic response of the vortex depends on these two parameters along with disk aspect ratio and defects in the disks.

A theoretical description of the eigenexcitation of a vortex state was provided by Guslienko et al. [207]. The translation mode of the vortex core has a circular trajectory in the absence of damping while it is spiral when damping was taken into account. The eigenfrequency of the vortex state depends on the dot aspect ratio \(\beta\) (thickness/radius) and increases linearly with increasing \(\beta\). This was experimentally shown by Novosad et al. [208].
Park and coworkers [143] have shown the dynamics of an isolated magnetic vortex in a permalloy disk with 50 Oe and 150 ps field pulse. Figure 1.24 shows the time evolution of magnetization for various sizes of the disks. They have attributed the low-frequency mode to the gyrotropic motion of the vortex core and the high-frequency mode to the precession of magnetization around the local internal field. The low-frequency gyrotropic modes are extremely long lasting as compared to the closure domain structures which are discussed later in this section. Guslienko et al. [207] have shown the frequency of the gyrotropic mode to be

\[ \omega_0 = \gamma M_s \xi^2 / 2 \chi(0), \]

where \( \gamma \) is the gyromagnetic ratio and \( M_s \) is the saturation magnetization.

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**Figure 1.24** (a) Magnetic force microscope image (left) and schematic of the vortex structure (right) of a 500 nm disk. The bright spot at the center of the disk in the image is due to the large z-component of the magnetization. (b)–(d) Experimental (left) and simulated (right) time-domain polar Kerr signals for vortex structures of diameters 2 micron, 1 micron, and 500 nm near the center of each disk. The low-frequency signal that is particularly prominent in the case of the 500 nm disk is the gyrotropic mode discussed in the text. Reprinted with permission from Ref. [143]. Copyright 2003 by the American Physical Society.
Susceptibility $\chi$ and the parameter $\xi$ depend on the magnetization distribution in the displaced vortex state. Two different models were proposed to describe the observed frequencies. One of them assumes rigid movement of the entire vortex which requires free poles and the alternative model avoids edge poles [207]. Pole free model was found to be in agreement by Park et al.

Phase-sensitive Fourier transform of the TRSKM data was developed by Buess et al. [209] to investigate different spin-wave modes in a lithographically patterned disk structure. Axially symmetric and nonaxially symmetric modes were excited using nonuniform excitation pulse (out-of-plane). Authors have observed a gradient of the excitation pulse in the plane of the vortex when the ferromagnetic disk was placed at the center of omega-shaped microcoil. Tipping field pulse of strength $\sim 50$ Oe was created by sending a current pulse of 100 ps rise-time. Polar Kerr signal was measured in stroboscopic manner. This nonuniformity generates nonaxially symmetric modes in the disk structures with vortex ground state. Using Fourier transform the authors were able to identify all the eigenmodes and back-Fourier transform technique allowed them to visualize the response of individual modes in the time domain. Axially symmetric modes (Fig. 1.25) and nonaxially symmetric modes were shown in their report. Figure 1.25 shows amplitude and phase of different modes. They have also discussed circularly polarized modes which are propagating in nature.

The interaction between spin waves and a vortex has been demonstrated by Park and Crowell [210]. The study was performed in ferromagnetic disks with diameters between 500 nm to 5 $\mu$m. The authors have found a splitting of the spin-wave modes when the diameter of the structure is less than 2 $\mu$m. The amplitude of splitting is of the same order as the vortex gyrotropic frequency. The polarity of the vortex core determines the relative phases of the splitting modes. At larger structures one single degenerate mode was observed. Such broken degeneracy in smaller structures shows the coupling between low-frequency excitations of the domain structure and the spin waves.

Azimuthal and radial modes in a disk structure can be selectively excited depending on the direction of the excitation field pulse. Zhu and coworkers [211] have experimentally shown these modes in the case of two different excitation field configurations. The study was performed on a permalloy disk of 700 nm diameter and 30 nm thickness and the transient field was applied to orthogonal directions using photoconductive switch technique. The dynamics was probed using TRMOKE experiment. The disk in between
Figure 1.25 Fourier transform (FT) of the axially symmetric eigenmodes. FT at each location results in a Fourier spectrum with five resonances (width \(\approx 1\) GHz). In (a)–(c), Fourier amplitude (left) and Fourier phase (right) at resonance for modes with radial symmetry. Their frequencies are 2.8 GHz (a), 3.9 GHz (b), and 4.5 GHz (c). (a) has the largest spectral weight, that is, it is dominating the motion and is responsible for the overall periodicity. The left-hand side of the images reports the results of micromagnetic simulation; the right-hand side are the experimental data. Because of finite spatial resolution, the experimental images extend over the border of the disk. Reprinted with permission from Ref. [209]. Copyright 2004 by the American Physical Society.
the transmission lines experiences an out-of-plane excitation while the disk on top of the transmission line feels an in-plane excitation. High-frequency oscillations are superimposed to the low frequency in the case of in-plane excitation. For out-of-plane excitation periodic oscillation with gradually decreasing amplitude are observed. A strong peak at 9.3 GHz (uniform precession) and smaller peaks at 2.4 (edge mode) and 13 GHz (radial mode with $n=2$ obtained from their simulations) were observed in the case of out-of-plane excitation. Peaks at 0.4 (gyrotropic frequency), 6 and 7.25 (normal modes), 9.5, 10.7, 11.6, and 12.8 GHz were observed for parallel excitation (Fig. 1.26). The modes at 6.0 and 7.5 GHz are found to be travelling spin waves in the angular direction. These modes are angular modes with $|m|=1$, which are well matched to the antisymmetric torque of the transient field. Separation between these two modes was attributed to the difference for scattering with vortex and magnons with clockwise and anticlockwise modes. Zaspel and coworkers [212] have shown gyrotropic mode and

![Figure 1.26](image)

**Figure 1.26** Fourier transforms of the experimental time-domain ferromagnetic resonance curves. (a) The dashed line (open circles) shows the spectrum for perpendicular pulsed field excitation. (b) The solid line (filled circles) is from the in-plane pulsed field excitation. *Reprinted with permission from Ref. [211]. Copyright 2005 by the American Physical Society.*
higher frequency modes in a soft permalloy disk. Vortex–magnon interaction was taken into account to model the excitation spectra analytically.

Splitting of the azimuthal modes due to core gyration has been mentioned in the previous sections. What happens when the disks with vortices are coupled in an array? This question attracted theoretical interests [213]. An experimental investigation on this query was provided by Barman et al. [134]. The experiment was done using time-resolved Kerr microscopy on dot arrays with different edge-to-edge spacing (150 nm = EP1, 200 nm = EP2, and 270 nm = EP3). Vortex state in these dots was confirmed using static MOKE hysteresis loops.

A clear splitting of the low-frequency gyrotropic mode was found in EP1 (150 nm) case (see Fig. 1.27). Splitting can also be observed in EP2 (200 nm) sample by applying a bias field. The results were discussed in terms of interaction between the dots where core polarization plays a significant role. Chirality and relative position of the cores also were found to be the key to the magnetostatic coupling which is negligible when edge-to-edge separation is more than 200 nm. The core polarization can be modified with the bias field, which also affects the coupling between the dots and led to splitting when the cores have opposite polarities. Splitting of a nondegenerate vortex translational mode was observed by Buchanan et al. [214] at large driving fields.

![Figure 1.27](image)

**Figure 1.27** The FFT spectra of the time-resolved Kerr rotations are shown for permalloy disk pairs EP1, EP2, and EP3 at bias fields $H_b = 0$ Oe, $H_b = 100$ Oe, and $H_b = -100$ Oe. Reprinted with permission from Ref. [134]. Copyright 2010 by the Institute of Physics Publishing group.
The authors argued that two peaks were associated with two steady-state solutions that differ in their phase lag with respect to driving field.

Compton and coworkers [215] have shown the effect of pinning in vortex dynamics. They have argued that pinning at the nanoscale is responsible for the variation of the gyrotropic frequency at low amplitude excitation. No such variation was found at higher excitation field suggesting a depinned configuration. Authors also argued that the variation in the frequency at low rf excitation fields is associated with the interaction of vortex core with a single nanoscale defect and it is independent of disk diameter. A correlation between physical microstructure and magnetic microstructure was shown by Chen et al. [216] in soft ferromagnetic films. Interconnection was demonstrated in the form of vortex state in permalloy disks for various thicknesses. Authors have concluded that the pinning mechanism is dominated by surface roughness, which can be generalized to virtually any soft ferromagnetic films.

A phase diagram of the vortex dynamics in the linear and nonlinear regimes was shown by Chen et al. [217]. At higher excitation fields, deformation of the vortex leads to core reversal via creation and annihilation of vortex–antivortex pair. For excitation below 1 Oe, gyrotropic frequency increases with decreasing excitation amplitude. This enhancement has been attributed to pinning. The phase map in a parameter space of drive amplitude and frequency shows clear boundary between pinned and unpinned regimes. A hysteresis was observed while sweeping the rf field, which suggests metastable gyrotropic orbits in the above mentioned phase space.

Magnetic domain wall dynamics have got lot of attention due to its potential applications in magnetoelectronic devices. Domain walls are interfaces between opposing magnetization directions. Different types of logic operations were demonstrated in 2005 by Allwood et al. [18]. One of the key parameters for fast operation with a domain wall–based device is mobility ($\mu$), which is the rate of change of domain wall propagation velocity ($v$). Mobility varies with external magnetic field ($H$) and they are connected with a simple relation: $v = \mu H$ [218]. Mobility is defined as $\mu = \gamma A / \alpha$ where $\gamma$, $A$, and $\alpha$ are the gyromagnetic ratio, domain wall width, and Gilbert damping parameter, respectively. For soft ferromagnetic material like permalloy ($A \sim 20$ nm, $\alpha = 0.01$) $\mu \sim 30–40$ m/s/Oe. Although a very low mobility value was found by Ono et al. [219] and that was attributed to the increased viscous damping of the domain wall motion due to edge roughness. Later in 2003, Atkinson et al. [220] reported the mobility to be $\sim 30$ m/s/Oe. The velocity of domain wall propagation was estimated by the authors in permalloy nanowire. The velocity was found to be
dependent on the applied magnetic field. Domain wall velocities of 500 m/s and 1500 m/s were reported for 40 Oe and 49 Oe, respectively. The same authors also argued that the lateral confinement does not affect the domain wall propagation velocity. A nonlinear domain wall velocity response with external magnetic field was later shown by Beach et al. [218]. They have introduced a breakdown field below and above which linear behavior was found with significantly different slopes.

Park et al. [143] have used optically triggered in-plane field pulse to investigate spin–wave dynamics of a closure domain structure. Different modes from confined geometries or localized domain wall were identified by means of space and TRMOKE technique. The measurements were performed in the remnant state which is closure domain structure with four quadrant separated by 90° Neel walls. The excitation geometry is shown in Fig. 1.28. Top and bottom quadrant experience a torque from the pulsed field in this configuration and they are opposite in sign. Temporal responses from these regions are shown in Fig. 1.28b and d. Figure 1.28c shows the dynamics in the domain wall region as marked by a dot in the inset. Low-frequency (0.8 GHz) oscillation was found in domain wall region as compared to high-frequency (1.8 GHz) mode in top and bottom quadrants. These two modes were further established with frequency domain images by calculating the Fourier transform of the time-domain data at each position. The low-frequency oscillation of the domain walls is remarkably different from an extremely long-lived gyrotrropic motion of the vortex core which was discussed earlier in this section.

9. TIME-DOMAIN COHERENT CONTROL OF PRECESSIONAL DYNAMICS

For the next generation data storage and memory devices, the data read–write time is an important factor and fast switching of magnetic bits has become an important topic of research since late 1990s [221]. However, continued precession of magnetization vector (ringing) after switching poses a limiting factor on the switching time. Materials with high damping may reduce the ringing time but cannot completely cut it down immediately after switching. Hence, coherent control of magnetization precession by tailoring magnetic field pulse shaping has started to attract the magnetics community from early 2000. It is believed that the use of coherent–control methods will allow the operation of magnetic data storage devices at fundamental frequency limits by eliminating the adverse effects of precessional oscillations.
9.1. Coherent Suppression of Precession in Magnetic Thin Films

In 2000, Bauer et al. [124,222] attempted such a coherent suppression of precession in 1.5 μm thick Lu$_{2.04}$Bi$_{0.96}$Fe$_5$O$_{12}$(BIG) single-crystal film onto (111)-oriented gadolinium gallium garnet (GGG). The film has a saturation...
magnetization $4\pi M_s = 1750 \text{ G}$, and a very low damping, leading to a strong ringing of magnetization after field pulse termination. In addition, the anisotropy fields and the saturation magnetization are small, and this results in low precession frequencies, which also cause a pronounced ringing. The above facts allowed to test a robust method of coherent suppression and an easy detection. The sample was saturated with an in–plane 4 Oe bias field to create a preferred, single-domain magnetization state with the direction of magnetization perpendicular to the plane of incidence. The dynamics was excited by an in–plane magnetic field pulse of 0.6 Oe (maximum tilt angle of 9° from the equilibrium) from a microstrip line applied perpendicular to the static field and the dynamics was probed by measuring the time–resolved Kerr rotation.

As a function of pulse duration ($T_{\text{pulse}}$), the amplitude of ringing after the pulse is terminated exhibits a full suppression at $T_{\text{pulse}} = 3.2, 6.4, 10.8, 13.8$, and $17.8 \text{ ns}$ with a pronounced ringing at intervening values of $T_{\text{pulse}}$ (see Fig. 1.29). The suppression is understood by using a single macrospin model,

![Figure 1.29](image)

**Figure 1.29** Time evolution of the $x$-component of the magnetization $M_x$ measured during and after field pulse excitation for different pulse durations $T_{\text{pulse}}$ as indicated. The beginning of the pulse launch is at $t = 0 \text{ ns}$. All measurements are performed in the center of the sample. *Reprinted with permission from Ref. [124]. Copyright 2000 by the American Institute of Physics.*
in which the field pulse must terminate at the moment at which most of the total magnetic energy is contained in the Zeeman energy of the pulsed field, which will be reduced to zero upon field termination. Inclusion of shape anisotropy and other energy terms would not affect the basic concept but would make the magnetization trajectory more complicated. The authors also took scanned images of the time-evolution of magnetization of the whole film showing a spatially uniform suppression. Later in 2001, the authors reported the coherent suppression in the same sample by applying a pulse with $T_{\text{pulse}} = 18$ ns and with rise and fall times of 0.4 ns and by varying the bias magnetic field [223]. A number of static bias fields satisfied the condition of coherent suppression. FFT of the ringing data revealed a number of spin-wave modes in the spectra, and the degree of suppression of two dominant modes as a function of the bias magnetic field (up to $\sim 40$ Oe) was reported. It was found that both modes are more efficiently suppressed at higher static magnetic field.

During the same time Crawford et al. [224] demonstrated a coherent suppression of a 50 nm thick permalloy film by a staggered step excitation. The staggered step was generated by sending a current step pulse in a shorted waveguide and the forward propagating and reflected step pulses at any given position of the waveguide forms two equal steps with a time delay given by $t_s = 2L/v_g$, where $L$ is the relative distance to the shorted end and $v_g$ is the group velocity of the electromagnetic wave in the transmission line (see Fig. 1.30). The time-resolved magnetization at different

![Diagrammatic sketch of the experimental arrangement used to demonstrate coherent control of magnetization dynamics in a thin-film sample. Reprinted with permission from Ref. [224]. Copyright 2000 by the American Institute of Physics.](image-url)
observations points along the waveguide was measured by second harmonic MOKE. Detection was in a coincidence detection mode, where the measured contrast is proportional to the $\gamma$-component of the magnetization. The observation points were selected by moving the waveguide with respect to the incident light. A complete suppression was achieved at $L=23\text{mm}$, which corresponds to a time delay of $380\text{ps}$. For materials with non-negligible damping, the second step excitation must be attenuated in proportion to the degree of damping that occurs during the first half cycle of magnetic response.

9.2. Coherent Suppression of Precession in Microscale Thin-Film Element

A coherent control of precessional switching in continuous thin films demonstrates the principle. However, for application of coherent precessional switching in magnetic storage and memory devices, achievement and understanding of coherent suppression in confined magnetic elements are desirable. In 2002, Schumacher et al. [225] reported the coherent suppression of precession in a $1 \times 4\mu\text{m}$ wide Ta 65 Å/NiFe 40 Å/MnIr 80 Å/CoFe 43 Å/Cu 24 Å/CoFe 20 Å/NiFe 30 Å/Ta 8 Å spin valve elements in a buried pulse line configuration. Magnetic field pulses ($H_p$) are created by sending voltage pulses (5 V maximum, pulse length 0.2–4 ns, rise-time 60 ps) through the pulse line. The transmitted pulse is recorded in a 50 GHz sampling oscilloscope. To detect the high-frequency GMR response dc currents of ±1 mA are applied through the spin valve via a bias tee. The voltage response pulse of the spin valve is picked up on one side of the sense line by the second oscilloscope channel while the other sense line contact is terminated to 50 V. Subtraction of two oscilloscope traces for positive and negative dc bias gives the spin valve resistance change from the current-independent crosstalk between the crossed lines. The precession period is first determined by allowing ringing to occur at an external static hard axis field of $H_s=91\text{Oe}$. The pulsed field length was then adjusted to slightly above the full precession period to suppress the precession. After a single precession cycle free layer magnetization ($M$) is still aligned with $H_s$. Only a small tilt between $M$ and $H_{\text{tot}}$ ($H_s + H_p$) is present when the pulse is switched off. As a consequence, almost no relaxation is needed to reach the final equilibrium magnetization parallel to $H_s$ and the ringing is suppressed. For two other pulses ($T_{\text{pulse}}=230$ and 760 ps), the pulse ends after ~0.5 and ~1.6 precession period, respectively. At these times the tilting angles between $M$ and the final precession axis $H_s$ are close to their maximum, leading to an increase in
precession amplitude and thus to strong ringing after pulse decay. Later Schumacher et al. [226] showed evidence of multiple coherent periodic precessional switching in a similar microscopic spin valve by ultrashort transverse field pulses as short as 140 ps. At high fields a phase coherent reversal is found revealing periodic transitions from switching to non-switching under variation of pulse parameters. At low fields a novel Gilbert damping-dominated reversal regime was found allowing highly efficient switching by fields below the hard axis anisotropy field, that is, below the quasistatic switching threshold.

In 2002, Gerrits et al. [227, 228] reported coherent control of precessional switching in a microscale elliptical permalloy elements by pulse shaping using a two pulse technique described in more details in Section 10. In 2005, Barman et al. [229] investigated spatial uniformity of the coherent suppression of precessional dynamics in microscale permalloy elements with 10 μm, width/thickness ratio: (S1, AR = 454), 7 μm (S2, AR = 318), and 5 μm (S3, AR = 227) width, all with thickness 22 nm and another 10 μm wide element with thickness 150 nm (S4, AR = 67). An oscillatory-shaped pulsed field \(h\) with 400 ps period was obtained due to impedance mismatches between the photoconductive switch, transmission line, and power supply. The precession is suppressed after one cycle of precession for specific bias field values, while ringing is observed for all other bias fields. For suppression of precession, the magnetization \(M\) must first become aligned with the total effective field and the torque upon the magnetization vanishes, and thereafter the total effective field must only change slowly. Since \(H\) and \(h\) lie within the sample plane, uniaxial anisotropy is negligible, and the demagnetizing field is antiparallel to the out-of-plane magnetization component, the magnetization can only become parallel to the total effective field if the magnetization lies within the sample plane. Suppression may then only be achieved for certain values of \(H\) such that \(H + h\) is parallel to \(M\) at this instant in time. Time-resolved scanning Kerr images were taken at specific time delays for all four samples to understand the spatial uniformity of the coherent suppression. Figure 1.31 clearly shows in black–gray–white scale that for S1, at bias fields where no coherent suppression occurs, the dynamics images show alternative bright and dark contrasts at different antinodes of the time-resolved data, whereas for \(H = 400\) Oe, where coherent suppression occurs after one cycle the image becomes gray, showing that magnetization has reached equilibrium. On the other hand for S3, at the bias field where no coherent suppression occurs, the images showed nonuniform precession near the edges parallel to the bias field. At \(H = 380\) Oe, where
coherent suppression seemed to occur, the images showed nonuniform contrast showing a spatially nonuniform coherent suppression in this sample. Further experimental works showed increased spatial nonuniformity in the time-resolved Kerr images for samples with reduced aspect ratio, suggesting that a simple pulse shaping scheme will not work and a more involved pulse shaping is required in confined magnetic elements.

9.3. Coherent Suppression in Magnetic Microstrips

In 2009, Barman et al. [230,231] experimentally demonstrated coherent suppression of precessional motion in a series of 100μm long permalloy microstrips of 5, 10, and 12μm width, in presence of multiple spin-wave modes. The lateral confinement of the microstrips causes spin-wave modes of frequencies adjacent to each other. Consequently, a position dependence of the suppression time is observed, which is consistent with the slight variation in the precession frequency with position along the short axis of the

Figure 1.31 The time-resolved Kerr images are shown for (a) S1, \( H = 145 \text{ Oe} \); (b) S1, \( H = 400 \text{ Oe} \); (c) S1, \( H = 732 \text{ Oe} \); (d) S3, \( H = 145 \text{ Oe} \); and (e) S3, \( H = 380 \text{ Oe} \). The delay times are shown under each image. Intensity images are presented in the column on the left for comparison. Reprinted with permission from Ref. [229]. Copyright 2005 by the American Institute of Physics.
microstrip. The frequency of the spin-wave modes along the short axis of the microstrip depends on the wave vector and hence on the width of the strip and, consequently, different suppression times for strips of different widths are observed. In fact the pulse duration ($\Delta T_p$) required for suppression reduces with the reduction of strip width as a result of the increase in the frequency of the confined modes. For $\Delta T_p$ values lying above and below the suppression time, ringing occurs with varying amplitude, which is also clear from the FFT spectra. Time-dependent micromagnetic simulations reproduced the suppression behavior. Simulated time-resolved images show that for an undershoot or overshoot of $\Delta T_p$ the magnetization does not become parallel to the effective field during the falling edge of the pulse and attains a new and an unstable precessional state. However, for a specific range of $\Delta T_p$ centered at around 500ps, the above condition is established and the precession stops coherently.

10. TIME-DOMAIN PRECESSIONAL SWITCHING

Reversal of magnetization at ultrafast time scale is required for faster performance needed in future memory and storage devices. Precessional switching has been shown as one of the promising techniques for this purpose. Using this technique the maximum time of reversal is half the precessional period, which is much faster compared to the domain wall motion (of the order of ns). To investigate this ultrafast dynamics it requires ultrashort pulses, which is a challenge to generate. Ultrafast response of the magnetization dynamics at the picosecond time scale was realized with the development of femtosecond laser sources, which were used to create ultrashort excitation. Subpicosecond dynamics have been observed experimentally, which show exciting results in the non-equilibrium state. Another issue in this reversal technique is known as ringing (precession after reversal), which has to be suppressed at the right time (half of the precession period) [124,232]. Some of the pioneering works in the area of precessional switching are discussed in the following paragraphs in this section.

Magnetization reversal via domain wall rotation was demonstrated by Choi et al. [233,234] with stroboscopic scanning Kerr microscopy. Experiments were performed in a $10 \times 2 \mu\text{m}$ permalloy strip placed on a transmission line. Bias field was applied along easy axis of the sample. A switching pulsed field was applied antiparallel to the bias field. In addition to these fields a transverse in–plane bias field was also used to manipulate the switching time
and reversal process. A remarkably faster reversal was observed in the presence of transverse bias field \( (H_t) \). The effect of \( H_t \) was clearly shown by the spatiotemporal domain images during reversal (see Fig. 1.32). Reversal mechanism was found to be dominated by the domain nucleation in the absence of transverse field whereas it was due to domain wall motion in the presence of the transverse field and the reversal was faster in the later case. The reversal time was reported to be 5 and 1.2 ns in the absence and presence of \( H_t \), respectively.

Precessional reversal with sub-ns field pulse along hard axis was shown by Schumacher et al. [235] by compensating the easy axis offset field. Switching was experimentally investigated by measuring the GMR of the device in crossed coplanar waveguide geometry. Switching was found for 360 ps field pulse with an amplitude of 170 Oe. The authors had shown precessional reversal using current pulse in a later report [236]. Current pulse of 120 ps duration and \( 5 \times 10^{11} \) A/m² current density was used to excite precessional dynamics in a spin valve stack. GMR response due to free layer reversal was measured and 2% GMR change was observed for full reversal for the above said current pulse. Precessional switching was confirmed by showing periodic variation of switching and nonswitching when the current pulse duration \( (T_{\text{pulse}}) \) was varied over a wide range (see Fig. 1.33). Switching occurs

Figure 1.32 Spatial profile of the \( M_x \) component as a function of time (ns) after the magnetic pulse was applied. \( H_l = 4.8 \) kA/m, while the transverse field is varied \( H_t = 0 \) and \( 5.2 \) kA/m. The field of view of each frame is \( 12 \times 4 \) μm and contains the entire \( 10 \times 2 \) μm sample. Reprinted with permission from Ref. [233]. Copyright 2001 by the American Physical Society.
when $T_{\text{pulse}} = (n + 1/2)T_{\text{precession}}$, where $T_{\text{precession}}$ is the precession period and $n$ is an integer. Back et al. [237,238] have shown that a precessional reversal is possible in Co/Pt thin film with a field pulse as short as 2 ps and amplitude of 184 kA/m when it is applied transverse to the magnetization.

There are several groups who have demonstrated the precessional switching in different ways. Lederman et al. [239] have shown the spontaneous thermal switching in single-domain Fe$_2$O$_3$ particles. Ju et al. [240] have shown the reversal of magnetization using subpicosecond laser excitation by large modulation of unidirectional exchange bias field across the ferromagnetic/antiferromagnetic bilayer films (NiFe/NiO). Koopmans et al. [102] have shown the dynamics in Ni thin films after excitation with femtosecond laser pulse. Pulsed electron beam from Stanford linear accelerator was used by Siegmann et al. [241] to generate ultrashort magnetic field pulse of 6 ps duration on CoPt films. This technique needs special requirement of linear accelerator, which hinders its general applicability. Doyle et al. [242]
and Hiebert et al. [119] have shown the generation of subnanosecond magnetic field pulse using microstrip line. Suppression of precession after the termination of the ultrashort field pulse is required for maximizing the switching speed. This was achieved by several groups by modulating the field pulse.

Gerrits and coworkers [228] have addressed the issue of ringing after precessional switching by engineering the field pulse profile. Two GaAs photoswitches were triggered independently by femtosecond laser excitation. Desired profile of the field pulse in time was obtained by superimposing these two pulses (see Fig. 1.34). Pump pulse 1 was used for initiating the precession and the pump pulse 2 acted as a stop pulse. The time delay between these two pulses can be precisely varied by using a delay line. Coherent suppression of the ringing was observed when this delay time was half the period of the precession. The reversal time was observed to

![Figure 1.34](image-url)  
**Figure 1.34** The schematic experimental setup. A sequence of two pump-laser pulses, which are separated by a distinct time delay, excites two GaAs-photoconductive switches. The generated current pulses are superimposed to produce one short and square-like magnetic field pulse. The field pulse is launched down a coplanar waveguide structure and excites the thin-film permalloy (Ni$_{81}$Fe$_{19}$) element at the end of the tapering. The inset is a microscopic photograph of the 8-nm-thin permalloy magnetic element that has an elliptical shape with dimensions of $8 \times 16 \mu m^2$. The vector- and time-resolved element response is measured by magnetization-induced second harmonic generation (MSHG) and the polar magneto-optical Kerr effect (MOKE). A high-reflectance infrared mirror (HRIR) is used to split the fundamental and second harmonic part of the beam. A photomultiplier tube (PMT) is used to detect the second harmonic photons. Reprinted with permission from Ref. [228]. Copyright 2002 by the Nature Publishing Group.
be 200 ps, which suggests a switching rate of 5 GHz for this particular experimental condition. In Fig. 1.35 the authors had shown a suppression of ringing with a stop pulse in contrast to the case where no stop was applied. A spin reorientation in few picoseconds was observed in antiferromagnetic TmFeO₃ material, which is faster than in the ferromagnetic materials [243].

Spatio-temporally and vectorially resolved vector Kerr microscopy was used by Hiebert et al. [244] to demonstrate precessional-based magnetization dynamics in thin-film magnetic elements. A scanning optical microscope was used to stroboscopically image the nonequilibrium magnetization states. The signal was phase sensitively detected. The time and spatial resolution was 50 ps and 0.65 µm, respectively. A large angle precessional dynamics was observed when a bipolar, square, in-plane field pulse to stadium-shaped permalloy element was applied. A static magnetic field (100 Oe) was applied along + x axis and a field pulse of −160 Oe (along −x axis) was created by the transmission line underneath the sample. The magnetic field changes from $H(t_0) = -60$ Oe to $H(t_f) = +100$ Oe. The long axis of the sample is along γ-direction. The shape of the sample defines an extra magnetic field along the long axis. When the field changes from −60 Oe to +100 Oe, magnetization feels a torque along γ-direction and starts fairly uniform reversal process. The torque drives the magnetization down out of the x–γ plane, which

![Figure 1.35 Switching by large-field excitation and suppression of ringing. Without a stop pulse, the system switches back to its initial state (open circles). After sending the stop pulse, the suppression of the ringing of the magnetization can clearly be observed (solid circles). The lines are guides to the eye. The low signal-to-noise ratio in the $M_x$ component results from the very weak longitudinal MSHG signal with an incoming polarization parallel and second harmonic polarization perpendicular to the plane of incidence ($p_{in} - s_{out}$). Reprinted with permission from Ref. [228]. Copyright 2002 by the Nature Publishing Group.](image)
in turn generates a demagnetizing field. Magnetization shows a damped resonant oscillation around the $+x$ axis with decreasing demagnetizing field. High-frequency spin-wave generation was ascribed to the reduction of the magnetization vector length. A small angle precessional dynamics was also studied in the same report in a permalloy disk, which was placed at the center of an omega-shaped loop to create an out-of-plane excitation field. Modal ferromagnetic oscillation was observed in the later case.

Precessional switching in a spin valve was reported by Kaka et al. [245] using magnetoresistance measurement technique. A bias field was applied along the long axis of the sample and the pulsed field was applied along the hard axis of the sample. The free layer consists of Ni$_{80}$Fe$_{20}$(2.5 nm)/Co(1 nm) and it is pinned along the easy axis. An antiparallel alignment of the spins in the free layer with the pinned layer causes higher resistance and parallel alignment has lower resistance. A pulsed field of 216 Oe amplitude and 230–325 ps duration was applied transverse to the bias and easy magnetization direction. Longitudinal bias field was used to negate the magnetostatic coupling field. Authors had shown that duration ($\Delta t$) and the amplitude ($H$) of the pulsed field is crucial for reversal via large angle precession. The energy ($\varepsilon$) required for the writing process is governed by the following relation: $\varepsilon; H^2 \Delta t$. Authors had shown that the probability of switching decreases with the pulse duration (Fig. 1.36).

![Figure 1.36](image)

**Figure 1.36** Probability of switching for a $0.45 \times 1.15$ $\mu$m device versus pulse duration. Dots indicate probability of complete switching and triangles indicate probability of remanent intermediate state. Inset: field pulses $H_p$ are along the hard axis. $H_L$ is the effective bias field in the easy-axis direction. $M$ gives the initial free layer magnetization direction. Reprinted with permission from Ref. [245]. Copyright 2002 by the American Institute of Physics.
11. SUMMARY AND FUTURE DIRECTIONS

The research in ultrafast magnetization dynamics has attracted a lot of attention due to its potential applications in magnetoelectronic devices and technology. Development of modern fabrication facilities down to sub-50nm has opened up new possibilities of applications like spin-transfer torque devices and magnonic crystals. Ultrafast optical techniques have made it possible to study the dynamics in subpicosecond time scale. Among various techniques TRMOKE has emerged as a very powerful technique for studying dynamics down to femtosecond time scale. Here, we have reviewed time-domain study of magnetization dynamics using TRMOKE technique in magnetic films and patterned magnetic structures. A theoretical background of magnetization dynamics has been provided for better understanding the experimental observations. In this respect LLG equation, Kittel mode, and dipole- and exchange-dominated spin waves are discussed in different wave vector and magnetic field configurations. The physical origin of the change of ellipticity in MOKE is explained. Later, the progress of developing TRMOKE technique since 1990 and the pioneering works in this area are highlighted. Ultrashort stimuli (optical, field, or current) are used to excite the dynamics down to femtosecond time scale and subsequently the response is recorded in stroboscopic manner with TRMOKE technique. Interest in the confined magnetic structures has increased significantly as we are approaching the technological limit in terms speed and performance. TRMOKE is an ideal candidate to probe the dynamics in the magnetic nanostructures. Nonuniform precessional dynamics in patterned structures was experimentally shown in 1997. In the year 2000 Freeman et al. demonstrated a method for imaging noise by time-resolved Kerr microscopy. Localized spin-wave modes were observed in magnetic microwires. Dependence of precessional frequency and damping on configurational anisotropy have been experimentally shown by Barman et al. The precession frequency showed a strong fourfold anisotropy superposed on a weak uniaxial anisotropy with the bias field orientation angle in a square permalloy element. Time- and vector-resolved MOKE technique has found direct application at the wafer level of magnetic hard disk writers. The experiment by Gangmei et al. provides an experimental confirmation of the well-known “flux beaming” mechanism in the yoke of a hard disk writer. Magnetic multilayers have been studied extensively due to their novel properties like spin-torque effect, GMR, spin valve, TMR, exchange spring magnets, PMA, exchange bias, etc.
Magnetization dynamics have been investigated in different multilayer systems to understand these properties in picoseconds time scale. Significant works in the multilayer systems have been discussed in this review. Size-dependent crossover of the precessional frequencies was observed in nanodot arrays in the 2005. When the nanoelements are very closely placed in an array, they can talk to each other and that gives rise to collective modes of the array. These modes are different from those arise from an isolated nanoelement. Some of the works in this direction are discussed in this review. The concept of magnonic crystal has been shown in this respect. One of the interesting findings is the dependence of the magnonic spectra on the lattice symmetry by Saha et al., which is explained in some detail here. TRMOKE was also employed greatly to understand vortex dynamics and domain wall dynamics. Femto and picosecond magnetization dynamics in single nanomagnets well below diffraction limit was studied by several groups. The works by Barman et al. employing a cavity-enhanced TRMOKE technique showed interesting variation of precession frequency and damping with the variation in nanomagnet size down to 125 nm. Later, in 2011 Rana et al. showed the precessional dynamics of magnetostatically isolated 50 nm permalloy dots are dominated by edge mode, where the center mode is suppressed. Precessional switching and coherent suppression of the precessional dynamics have been studied using TRMOKE technique and pioneering research works in this field have been explained.

Applications of TRMOKEs to study the dynamics of magnetic materials and metamaterials were not limited to those presented above. There have been a large numbers of important and pioneering studies in other systems such as metallic and alloy thin films, in particular Fe–Pt thin films [246,247], CoFeB [248], Mn–Ga [249], and CoCrPt [250] thin films with high PMA, CoCrPt:SiO$_2$ granular media [251] and Co$_2$MnAl [252] and Co$_2$MnSi [253] full Heusler alloy films. A sizeable amount of works on the spin, carrier, and phonon dynamics on half-metallic oxide thin films such as CrO$_2$ [254,255], La–Ca–Mn–O [256–258] and La–Sr–Mn–O [259] films have also been reported. In 2009, Müller et al. [114] reported ultrafast demagnetization in various half-metal oxides, Heusler alloys, Fe$_3$O$_4$ and Ni thin films by optical pump–probe technique and found a correlation between the spin polarization of the material with the demagnetization time. A strong effort has been made by primarily the Nijmegen group on optomagnetic switching of Tm–Fe–O and Gd–Fe–Co systems by using strong temperature dependence of anisotropy or in a nonthermal way by using inverse Faraday effect [243,260–262] leading toward possible applications in all-optical magnetic recording in femtoseconds time scale.
In addition to various thin films and patterned micro- and nanoelements, some interesting studies of the localized and extended spin-wave modes in antidots on magnetic thin films have started to appear in the literature [263–266]. Studies of ultrafast demagnetization, relaxation and magnetization precession of magnetic nanoparticles synthesized by chemical routes and ion implantation methods have also been reported [267–271]. Some of these nanocrystals reveal unique properties such as high resonance frequency, strong effective damping, and electrically insulating character, which are favorable for applications in sensors.

Time-resolved magneto-optical Kerr microscopy has been proven to be a very unique and powerful technique to study ultrafast dynamics in different types of magnetic systems. Given the progress in last two decades in terms of scientific knowledge and technological development, it would be of interest to extend the research for faster operation and efficient control of the magnetoelectronic devices. Devices like magnonic crystals, STNOs, and nanowire waveguides are very crucial for spin-based electronics, and a detail understanding of propagating and standing spin-wave modes using TRMOKE technique will be crucial. Spintronic devices also face challenges from the quality of fabrication down to sub-50 nm length scale and bulk fabrication limitations. A parallel technological development of these fabrication issues will also be necessary.

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