Field-Induced Magnetization in Nanostructures

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Abstract The field-induced impact on magnetic nanostructures with a large spin-orbit interaction, consisting in magnetization reversal under ultra-short circularly polarized laser pulses or unipolar electric field pulses are studied. Using the magneto-optical method and a pump-probe technique based on the Kerr and Faraday effects, we have established features and conditions of the magnetization reversal in magnetic nanostructures under femtosecond circularly polarized laser pulses. It is shown that mechanisms of such the laser-induced impact is a complex process of laser-induced thermal demagnetization of magnetic sublattices with subsequent biasing by internal magnetic fields of different nature. The interfacial voltage-controlled magnetic anisotropy in magnetic nanostructures is studied. In the framework of the model, based on the Stoner magnetization and the Rashba spin-orbit interaction the conditions of the electric control of the perpendicular magnetic anisotropy and the magnetization switching are considered.

Keywords: laser-induced and electric-induced remagnetization, spin-orbit interaction, magnetic nanostructures


1. Introduction

Physical limits of remagnetization speed and energy consumption are one of fundamental problems of magnetism physics, which has a crucial significance for creation of systems of high-speed magnetic recording and readout of information. Growth of attention to this problem is related to modern achievements of nanotechnologies, by possibilities of production of new magnetic nanostructures with predetermined physical properties, and development of the short-time pulsed sources of laser radiation and electric field. The prospect of the solution of this problem is related to the use of the impact of short-time laser and electric field impulses on the magnetic, specifically ferrimagnetic, multilayered nanostructures including tunnel magnetic junctions [1], that can lead to magnetic state variations and the remagnetization effect.

The key role belongs to the effects connected with the laser-controlled and electric field-controlled spin manipulation including the spin-polarized electron transport in magnetic nanostructures that constitute the subject of the modern magnetism physics, spintronics and microelectronics. The corresponding field-induced magnetic transitions in magnetic nanostructures are determined both by their spin and electron structure and by physical characteristics of the external fields. Corresponding magnetic materials must possess by high enough values of magnetic anisotropy, spin-orbit interaction, magneto-optical susceptibility and the spin polarization of conduction electrons.

The directed fast-speed laser-induced magnetic impacts is realized via the laser-induced thermal demagnetization with subsequent bias by effective internal magnetic fields of different nature subject to the magnetic structure of materials and laser radiation characteristics. The effective internal magnetic bias fields can be related to the inverse magneto-optical Faraday effect (see [1,2]) and to the effects of the transient ferromagnetic-like state in combination with the relaxation of the exchange antiferromagnetic interaction under the pulsed laser-induced thermal demagnetization [3,4]. In the first case, the bias magnetic field \( H_F \) arises only under circularly polarized laser irradiation. In the second case, the bias field \( H_{TR} \) is the internal transient field caused only by the different speeds of the laser-induced thermal demagnetization and it is independent on the radiation polarization.

The indirect laser-induced magnetic impact is realized via the laser-induced spin-polarized electron current, caused by the laser-induced kinetic effect of a spin-polarized electron current associated to momentum-selective excitations and a photon pressure [5]. In so doing, the laser magnetic control is realized via the exchange interaction between the spin-polarized electron current and localized magnetic moments of a magnetic nanostructure [6].

The electric field-induced magnetization occurs via the Rashba spin-orbit interaction [7] assuming a two-dimensional electron system with the broken spatial inversion symmetry. The pulsed electric field exerts the effect of spin splitting of the electronic band spectrum with the corresponding spin polarization of surface states of a metal magnetic nanolayer. The effective bias
magnetic field, in this case, corresponds to the combination of the Rashba spin-orbit interaction and the s-d-exchange interaction between conduction electrons and the localized of magnetic nanolayers.

The paper is organized as follow. In section 2, the dynamics of magnetization switching in tunnel magnetic nanostructures is studied. In section 3, features of the passage of the spin-polarized current through tunnel magnetic nanostructures are established. In section 4, the model description of the laser-induced remagnetization including a thermal demagnetization is considered. The magnetic memory based on the interaction between the spin polarized current and localized magnetization in multilayer tunnel magnetic nanostructures is proposed. In section 5, the electric field-induced magnetization in magnetic nanostructures based on transition 3d magnetic including the Rashba spin-orbit interaction and the band structure. In this case, the remagnetization of the magnetic junction can be caused by the exchange s-d interaction of the laser-injected spin-polarized current with the localized lattice magnetic moments (see [6]). The effective internal magnetic field $H_{sd}$ of that interaction constitutes from two components, $H_{sd} = H_S - H_{inj}$. The first component $H_S$ is related to the s-d-interaction of the transverse component (with respect to the magnetic moment of the injected layer) of the magnetic moment of the spin-polarized current. The second component $H_{inj}$ is related to the s-d interaction of the laser-induced thermal demagnetization result in the remagnetization, which is accompanied by a tunneling magneto-optical effect and the results of the magnetooptical measurements are represented in Figure 2.1.

2. Dynamics of Magnetization Switching

The influence of femtosecond pulses of circularly polarized laser radiation on magnetic states and the conductance of the spin-polarized electron current was studied for the ferrimagnetic TbCoFe-based and the ferromagnetic CoFe-based nanostructures. Features of the laser-induced magnetization reversal were studied for magnetic nanostructures $\text{Al}_2\text{O}_3/\text{Tb}_2\text{Co}_{10}\text{Fe}_{70}/\text{Al}_2\text{O}_3$, $\text{Al}_2\text{O}_3/\text{Tb}_{25}\text{Co}_{5}\text{Fe}_{70}/\text{Al}_2\text{O}_3$, $\text{Al}_2\text{O}_3/\text{Tb}_{19}\text{Co}_{5}\text{Fe}_{76}/\text{Al}_2\text{O}_3$ and $\text{Al}_2\text{O}_3/\text{Co}_{30}\text{Fe}_{70}/\text{Al}_2\text{O}_3$ with a single magnetic nanolayer. The laser-induced magnetization reversal together with the laser-induced tunnel magneto-resistance (TMR) effect we studied for magnetic junctions $\text{Al}_2\text{O}_3/\text{Tb}_{22}\text{Co}_{5}\text{Fe}_{73}/\text{Pr}_6\text{O}_{11}/\text{Tb}_{19}\text{Co}_{5}\text{Fe}_{76}/\text{Al}_2\text{O}_3$ and $\text{Al}_2\text{O}_3/\text{Co}_{30}\text{Fe}_{70}/\text{Pr}_6\text{O}_{11}/\text{Co}_{30}\text{Fe}_{70}/\text{Al}_2\text{O}_3$ with the PrO-based isolating barrier nanolayer. The TbCoFe-based and CoFe-based layers are characterized by uniaxial out-plane and in-plane magnetic anisotropy, respectively. The magnetization pinning of nanolayers $\text{Tb}_{25}\text{Co}_{5}\text{Fe}_{70}$ and $\text{Co}_{30}\text{Fe}_{70}$ considerably larger than the nanolayers $\text{Tb}_{19}\text{Co}_{5}\text{Fe}_{76}$ and $\text{Co}_{30}\text{Fe}_{70}$. The laser-induced magnetization dynamics was studied with the help of the magneto-optical pump-probe method, based on the magneto-optical Kerr and Faraday effects. The Nd-YAG and He-Ne lasers generated the pumping pulses (with a duration $\tau_p \approx 130$ fs) and corresponding probe pulses, respectively.

The direct laser impact on the magnetization occurs via the interaction of the circularly polarized photons with spin-polarized electrons of a magnetic medium. The Raman-like photon excitation of the electrons together with a spin-orbital interaction are accompanied by the spin-flip and the remagnetization (see [2]) providing the inverse magneto-optical Faraday effect with the effective internal bias field ($H_F$). The pulsed laser irradiation causes heating and demagnetization that in combination with the laser-induced effective magnetic field can lead to variation of magnetic states and the remagnetization. The direct femtosecond laser-induced impacts can results in the internal effective bias magnetic field $H_{TR}$, connected with the transient ferromagnetic-like state, caused by different demagnetization rates of ferrimagnetic sublattices together with the exchange interaction relaxation [3,4]. However, in the considered case of TbCoFe-based ferrimagnetics this bias field $H_{TR}$ can play role of amplification of the effective bias field of the inverse magneto-optical Faraday effect (see [4]).

The indirect laser impact on magnetic states in nonuniform multilayered magnetic nanostructures occurs via the laser-induced spin-polarized electron current between magnetic nanolayers. In this case, the remagnetization of the magnetic junction can be caused by the exchange s-d interaction of the laser-injected spin-polarized current with the localized lattice magnetic moments (see [6]). The effective internal magnetic field $H_{sd}$ of that interaction constitutes from two components, $H_{sd} = H_S - H_{inj}$. The first component $H_S$ is related to the s-d-interaction of the transverse component (with respect to the magnetic moment of the injected layer) of the magnetic moment of the spin-polarized current. The second component $H_{inj}$ is related to the s-d interaction of the laser-induced longitudinal spin component (with respect to the magnetic moment of the injected layer) which is characterized by a nonequilibrium distribution [6]. The mentioned effective internal magnetic fields together with laser-induced thermal demagnetization result in the remagnetization, which is accompanied by a tunneling magneto-optical effect. The results of the magneto-optical measurements are represented in Figure 2.1.

![Figure 2.1](image-url) Figure 2.1. The time dynamics of the laser-induced remagnetization under the high-power polarized laser pulses. (I) Probe laser pulses pass through the nanolayers $\text{Al}_2\text{O}_3/\text{Tb}_{25}\text{Co}_{5}\text{Fe}_{70}/\text{Al}_2\text{O}_3$ (1) and $\text{Al}_2\text{O}_3/\text{Tb}_{19}\text{Co}_{5}\text{Fe}_{76}/\text{Al}_2\text{O}_3$ (2) irradiated by circularly polarized pulses. (II) The probe laser pulses are reflected from the nanolayer $\text{Tb}_{25}\text{Co}_{5}\text{Fe}_{70}$ of the magnetic junction $\text{Al}_2\text{O}_3/\text{Tb}_{25}\text{Co}_{5}\text{Fe}_{70}/\text{Pr}_6\text{O}_{11}/\text{Tb}_{19}\text{Co}_{5}\text{Fe}_{76}/\text{Al}_2\text{O}_3$ irradiated from the side of nanolayer $\text{Tb}_{25}\text{Co}_{5}\text{Fe}_{70}$ by circularly (1) and linearly (2) polarized pulses. (III) Probe pulses are reflected from the nanolayer $\text{Co}_{30}\text{Fe}_{70}$ of the tunnel magnetic junction $\text{Al}_2\text{O}_3/\text{Co}_{30}\text{Fe}_{70}/\text{Pr}_6\text{O}_{11}/\text{Co}_{30}\text{Fe}_{70}/\text{Al}_2\text{O}_3$ irradiated from the side of the nanolayer $\text{Co}_{30}\text{Fe}_{70}$ by circularly polarized laser pulses. Arrows denote magnetization directions in adjacent nanolayers.

As it can be seen from the temporal behavior of curves I (Figure 2.1), the dynamics of the laser-induced magnetization for the infrared laser radiation is different for nanolayers $\text{Tb}_{19}\text{Co}_{5}\text{Fe}_{76}$ (the curve 1) and $\text{Tb}_{25}\text{Co}_{5}\text{Fe}_{70}$ (the curve 2). The dominant interaction of light with the iron magnetic sublattice results in that the nanolayers $\text{Tb}_{19}\text{Co}_{5}\text{Fe}_{76}$ more rapid changes its magnetization sign than the nanolayer $\text{Tb}_{25}\text{Co}_{5}\text{Fe}_{70}$ during the laser-induced remagnetization. For the ultraviolet laser radiation, when the interaction between light and rare-earth. Tb sublattice is dominant; the magnetization is defined by the Tb atom content. Then the remagnetization of nanolayer $\text{Tb}_{25}\text{Co}_{5}\text{Fe}_{70}$ will change itself sign rapid than the nanolayer $\text{Tb}_{19}\text{Co}_{5}\text{Fe}_{76}$ [4].
Due to curves II (Figure 2.1), the laser-induced magnetization reversal of the TbCoFe-based ferromagnetic junction under laser pulses occurs faster at the circular (the curve 1) than linear (the curve 1) polarization of laser radiation. Such a rate difference is caused by the grater total effective bias field in the case of the circular polarization \((H_T = H_F + H_{sd} + H_{TR})\) and the less total effective bias field in the case of the linear polarization \((H_T = H_{sd} + H_{TR})\). The laser-induced remagnetization of the CoFe-based tunnel ferromagnetic junction (the curve 111 in Figure 2.1) is related to the dominant effective bias field \(H_{sd}\) of the \(s\)-\(d\) exchange interaction between the laser-induced spin-polarized current and localized magnetic moments. The above-mentioned laser-induced remagnetization in the TbCoFe-based ferromagnetic and CoFe-based ferromagnetic junctions is accompanied by the tunnel magnetoresistance, i.e., by the dependence of the spin-polarized current through a magnetic junction on the magnetization configuration of its nanolayers.

3. Passing of the Spin-Polarized Current

The PrO barrier layer represents a large-gap semiconductor \([8]\) similarly to the known \([9,10,11]\) case of the MgO-based barrier in the magnetic junction Fe/MgO/Fe. The PrO-based tunnel barrier as well as the MgO-based tunnel barrier is characterized by a large enough tunnel transparency and, consequently, a large TMR effect under external magnetic field \([8]\).

The pulsed laser-induced remagnetization of magnetic nanolayers results in the pulsed change of the spin-polarized electron current through a tunnel magnetic junction. The laser-induced remagnetization of magnetic nanolayers results in the pulse change of the spin-polarized current through the tunnel magnetic junctions.

Such the change is related to the dependence of the density of states of the spin-polarized electrons on the magnetic moment of the nanolayers. The tunnel current through considered magnetic nanostructures in the most general form is expressed via the tunneling probability \((T)\) of the Bloch electrons from one magnetic nanolayer through the tunnel barrier nanolayer into an adjacent magnetic nanolayer.

The current density of the spin-polarized electrons from the magnetic nanolayer with the chemical potential \((\mu_1)\) into the magnetic nanolayer with the chemical potential \((\mu_2)\) is described by the expression \([10]\),

\[
J_{\mu_1 \rightarrow \mu_2} = \frac{e}{(2\pi)^2} \int d^3 k v_z(k) f(\mu_1) \sum_{k'} T(k, k') \tag{3.1}
\]

where \(e\) is an electron charge, \(v_z(k)\) is the velocity of electrons with the wave vector \(k\) along the axis \(z\), perpendicular to the plane of the barrier nanolayer, \(f(\mu_1)\) is the Fermi-Dirac distribution function, \(T(k, k')\) is the tunneling probability of electrons with the change of the electron wave vector from \(k\) on \(k'\). Performing the integral over \(k_z\) yields

\[
J_{\mu_1 \rightarrow \mu_2} = \frac{e}{2\pi} \frac{d^3 k v_z(k) f(\mu_1) \sum_{k'} T(k, k')}{k_z^2} \tag{3.2}
\]

where \(e = \varepsilon(k')\) is electron band spectrum. Taking into account \(3.2\), the net current, including opposed tunnel currents between magnetic nanolayers, can be expressed as

\[
I = \frac{e}{\hbar} \int d\varepsilon \sum_{k} T(k), \tag{3.3}
\]

where \(j\) is the number of the Bloch state for a given value of the transverse component \(k_j\) of the electron wave vector. This results in the Landauer formula

\[
G = \frac{e^2}{\hbar} \sum_{k} T(k), \tag{3.4}
\]

which describes the conductance of the magnetic tunnel nanostructures subject to features of the electron transmission through the tunnel nanolayer.

Calculation of the tunneling probability assumes allowing for the quantum state structure of spin-polarized electrons in the magnetic nanolayers, the structure of decaying evanescent electron states in the tunnel barrier nanolayer PrO and near interfaces. In the considered tunnel magnetic junctions the barrier nanolayer is wide-band semiconductor, in which similarly to the case of the tunnel magnetic nanostructure Fe/MgO/Fe, the tunneling conductance is realized via the two tunneling channels corresponding majority and minority spin polarized along and oppositely to the magnetic field.

The existence of the electron Bloch states with the relatively high spin polarization in the magnetic nanolayers and their overlap with slowly decaying quantum states in the barrier cause the high enough tunneling conductance in the majority channel. In the minority channel, the relatively large tunneling conductance related to interfacial resonance states caused by quantum interference between the decaying states in the barrier layer that provides efficient tunneling the Bloch electrons through the barrier. The mentioned large electron spin polarization together with the effective tunneling lead to the relatively large TMR effect owing to the corresponding change of an electron density state at the current transition between nanolayers with opposite magnetizations.

The explicit dependence of the spin-polarized current through the tunnel magnetic junction on the density of electron states, dependent on magnetic moments of the layers, follows from the formula (3.1). At given magnetizations of adjacent magnetic nanolayers of the magnetic junction, in the approximation of the mean tunneling probability \(T\), the components of spin-polarized electron current are described by the expression

\[
I_{\sigma, m, m'} = \frac{4e\pi^2}{\hbar} \int_{-\varepsilon_0}^{+\varepsilon_0} T g_1(\varepsilon - eV, \Sigma_1^{m}) g_2(\varepsilon, \Sigma_2^{m'}) \left[ f(\varepsilon - eV) - f(\varepsilon) \right] d\varepsilon, \sigma = (\uparrow, \downarrow), \tag{3.5}
\]
where \( g_{i \sigma}(\epsilon, \Sigma_i^m) \) \((i = 1, 2)\) is the electron state function in the \( i \)-th magnetic nanolayer with the magnetization \( \Sigma_i^m \), where the index \( m \) takes the values 1 and 2 if the electron spin is parallel and antiparallel to this magnetization, respectively. In addition, the function \( f(\epsilon) \) is the Fermi-Dirac distribution. The influence of the magnetization \( \Sigma_i^m \) adds up to the energy shift of the density of states function on the magnitude of the exchange interaction between the spin magnetic moments and the localized magnetic moment in nanolayers. As the density state function depends on a magnetization direction (as it is visible from (3.5)), the spin-polarized electron current is dependent on the magnetization configuration of the magnetic nanolayer.

At weak bias field and voltages \((V)\), less than the width of energy band spectrum of the Bloch electrons, the expression for the tunneling spin-polarized current can be represented in the form

\[
I_{\sigma,m,m'} = V \frac{4 \hbar^2 \alpha^2}{T} \epsilon \nonumber
\]

\[
\times \int_{-\infty}^{+\infty} g_{i \sigma}(\epsilon, \Sigma_i^m) g_{2 \sigma}(\epsilon, \Sigma_2^m) \frac{\partial f(\epsilon)}{\partial \epsilon} d\epsilon
\]

\[(3.6)\]

Hence, at low temperatures, where the distribution function \( f(\epsilon) \) takes on the form of the Heaviside step function, the following formula can be obtained for the tunneling conductance in majority and minority channels

\[
G_{\sigma,m,m'} = \frac{4 \hbar^2 \alpha^2}{T} \epsilon \nonumber
\]

\[
g_{1 \sigma, \Sigma_1^m} g_{2 \sigma, \Sigma_2^m}
\]

\[(3.7)\]

Here density-of-states functions are taken at the Fermi energy \( \epsilon_F \), i.e. \( g_{i \sigma, \Sigma_i^m} = g_{i \sigma, \Sigma_i}(\epsilon_F, \Sigma_i^m) \). Due to (3.7) at low temperatures and voltages, the components of the tunneling conductance are determined by density state functions of spin-polarized electrons at the Fermi energy in the magnetic layers.

At the parallel magnetization orientation of the adjacent magnetic nanolayers

\[
G_p = G_{11} + G_{22} \propto g_{1 \uparrow,1 \uparrow} g_{2 \uparrow,2 \uparrow} + g_{1 \downarrow,1 \downarrow} g_{2 \downarrow,2 \downarrow}
\]

\[(3.8)\]

At the antiparallel magnetization orientation (\( \Sigma_1 \parallel \Sigma_2 \))

\[
G_{AP} = G_{12} + G_{21} \propto g_{1 \uparrow,2 \downarrow} g_{2 \downarrow,1 \uparrow} + g_{1 \downarrow,2 \uparrow} g_{2 \uparrow,1 \downarrow}
\]

\[(3.9)\]

Due to equations (3.8) and (3.9), the tunnel magnetoresistance effect (TMR), i.e. the relative change of the tunneling conductance under switching the magnetization configuration of the magnetic junction between parallel and antiparallel states, is described by the expression

\[
TMR = \frac{G_p - G_{AP}}{G_{AP}} = \frac{2P_l P_l}{1 - P_l P_l},
\]

\[
P_l = \frac{g_{1 \uparrow,1 \uparrow} - g_{1 \downarrow,1 \downarrow}}{g_{1 \uparrow,1 \uparrow}}
\]

\[(3.10)\]

where it is allowed for the relations \( g_{1 \uparrow,2 \downarrow} = g_{1 \downarrow,2 \uparrow} \) and \( g_{1 \downarrow,2 \uparrow} = g_{1 \uparrow,2 \downarrow} \). As it is visible from (3.10), TMR effect is connected by the direct dependence with spin polarizations \( P_l \) of magnetic nanolayers.

Thus, the laser-induced change of conductivity of the tunnel magnetic junctions are caused by changes of density of states of spin-polarized electrons near a Fermi level because of the transition between parallel and antiparallel magnetization configuration of the junctions. Measured results of the conductivity dynamics of the tunnel magnetic junctions are represented in Figure 3.2.

**4. Modelling the Induced Magnetization**

The laser-induced magnetization dynamics of the explored tunnel magnetic junction, including the temperature dependence of the magnetization magnitude, in the approximation of effective magnetic moments of nanolayers can be described by the macroscopic Landau-Lifshitz-Bloch (LLB) equation [12,13,14],

\[
\frac{\partial m_i}{\partial t} = -\gamma \left[ m_i \times H_{eff} \right] + \frac{\gamma \alpha}{m_i} \left[ m_i \times \left( H_{eff} + \zeta \right) \right] m_i
\]

\[\text{(4.1)}\]

which was obtained in a mean-field approximation from the classical Fokker-Planck equation for individual spins interacting with a heat bath. In (4.1) \( \gamma = \gamma / (1 + \lambda^2) \), where \( \gamma \) is the gyromagnetic ratio and \( \lambda \) is a microscopic parameter that characterizes the coupling of the individual atomistic spins with the heat bath. It is visible from (4.1), that a spin polarization \( m_i \) has no constant length and is temperature dependent. The coefficients \( \alpha \) and \( \alpha \) are dimensionless longitudinal and transverse damping parameters. Thermal fluctuations are included as an additional noise terms \( \zeta_i(t) \) with \( t = (\lambda, \mu) \) and

\[
\langle \zeta_i(t) \rangle = 0
\]

\[
\langle \zeta_i(t) \zeta_j(t') \rangle = \frac{2k T \omega_i}{\gamma M_s \alpha} \delta(\omega_i \delta_i \delta_j \delta(t))
\]

\[\text{(4.2)}\]
where \( i, j \) denote lattice sites and \( \eta, \mu \) denote the Cartesian components. Here, \( \Delta^3 \) is the volume of the micromagnetic cell and \( M_0^3 \) is the value of the spontaneous magnetization at zero temperature. The damping parameters below and above the magnetic phase transition temperature \( T_M \) are described by the expressions \( \alpha_\parallel = 2T / \left( 3T_M \right), \alpha_\perp = \lambda (1 - T / T_M) \) and \( \alpha_\perp = 2\lambda T / \left( 3T_M \right), \) respectively.

The effective magnetic field can be written as \( H_{\text{eff}}^i = H_{1-\text{ind}} + H_{a-\text{ex}} \). Here the first ingredient \( H_{1-\text{ind}} \) is caused by the laser-induced electron excitations which includes both the effective internal magnetic field of the inverse magneto-optical Faraday effect \( (H_F) \) and the internal magnetic field \( (H_{\text{sd}}) \), generated by the \( s-d \) exchange interaction of laser-injected spin-polarized currents with localized spins of magnetic lattice in the spatially inhomogeneous magnetic junction. The second ingredient \( H_{a-\text{ex}} = H + H_{a} + H_{\text{ex}}^i \) is given by

\[
H_{a-\text{ex}} = \frac{1}{2\xi_\perp} \left( \frac{1 - m_\perp^2}{m_{\text{eq}}} \right) m_\parallel \theta (T_M - T)
+ \frac{1}{2\xi_\perp} \left( 1 + \frac{3T_m m_\perp^2}{5(T - T_M)} \right) m_\parallel \theta (T - T_M),
\]

(4.3)

where the susceptibility \( \xi_\perp = \partial m_\parallel / \partial H_F \). The anisotropy and exchange fields are given by

\[
H_a^i = -\left( \frac{m_\perp^2 e_x + m_\parallel e_y}{\xi_\perp} \right)
\]

and

\[
H_{\text{ex}}^i = -\frac{A}{m_\perp^2 M_0^3 \Delta^3} \sum_{j \in \text{neigh} (i)} \left( m_j - m_i \right),
\]

respectively. Within the context of the LLB equation, field components parallel to the local magnetic moment can change the length of the magnetization vector. In the limit, \( T \to 0 \) the longitudinal damping parameter \( \alpha_\parallel \) vanishes and with \( |m| = m_{\text{eq}}(0) \) the LLB equation goes over to the usual Landau–Lifshitz–Gilbert (LLG) equation [12]. The temperature dependent parameters in (4.1), i.e. longitudinal, transverse susceptibilities, and the temperature variation of the magnetization, \( \chi_\parallel (T) \), \( \chi_\perp (T) \) and \( m_\parallel (T) \) are determined using an Langevin dynamics combined with the LLG equation for each spin, i.e., by its stochastic modification

\[
S_i = \gamma \left[ S_i \times H + \lambda S_i \times S_i + \sqrt{\lambda} S_i \times H_i \right]
\]

(4.4)

where the internal field \( H_i = \partial H / \partial S_i + \xi_i (T) \). Thermal fluctuation of the mentioned parameters are included as an additional noise term in the internal field with \( < \xi_i (t) >= 0 \) and \( < \xi_i^k (t) \xi_i^l (t) >= 2\delta_{ij} \delta_{kl} \kappa B T \mu_S / \gamma \).

The systems (4.1) and (4.4) has solution for the magnetization \( m(t) \), which at the temperatures close to the magnetic phase transition temperature \( T_M \) tends to zero (that means a demagnetization process). Then at cooling, one tends to the magnitude with the sign opposite to initial value (that means the magnetic switching). The system (4.1) and (4.4) determine the conditions for parameters providing the magnetic switching. It turns out that the magnetic switching only occurs within a narrow range of parameters for the laser pulse. The realization of the magnetic switching assumes the suitable combinations of laser pulse duration and intensity.

The component \( H_{sd} \) of the effective magnetic field \( H_{\text{eff}} \) in (4.1) expresses via the \( s-d \)-exchange interaction \( U \) between the spin-polarized current and the lattice magnetization \( m \) as

\[
H_{sd} (x,t) = -\frac{\delta U_{sd}}{\delta m(x,t)}
= -\alpha \frac{\delta}{\delta m(x,t)} \int dx' m_{eq}(x',t)m(x,t)
\]

(4.5)

where \( m_{\text{eq}}(x,t) \) is the magnetization of the laser-induced spin-polarized current and \( L \) is thick of a magnetic nanolayer of the tunnel magnetic junction. Since the magnetization \( m = m(x,t) \) is connected with the magnetization flux density \( J \) by the continuity equation [6]

\[
\frac{\partial m_{el}}{\partial t} + \frac{\partial J}{\partial x} + \gamma \eta \left[ m_{el} \times m \right] - \frac{m_{el} - \overline{m}_{el}}{\tau} = 0,
\]

(4.6)

where \( \overline{m}_{el} \) is an averaged magnetization, \( \tau \) is a relaxation time with respect to a local equilibrium state, then the effective magnetic field \( H_{sd} = H_{sd}(J) \), i.e., it depends on the magnitude of the laser-induced current and the intensity of laser pulses.

The continuity condition of the magnetization flux near to the interface between continuity adjacent magnetic layers determines the boundary conditions for (4.6) that allows describing the magnetization dynamics under the laser-induced spin-polarized electron current. The continuity condition for the traverse components of the magnetization flux near to interface between adjacent magnetic layers result in a transfer of torque moment from labile electrons to lattice moments. The corresponding traverse component \( (H_{sd\perp}) \) of internal magnetic field \( H_{sd} \) can result in magnetic switching in a small region near to the interface at excess of threshold intensity of laser pulses [5]. Such the spin torque effect assumes spin dissipation.

At the same time, the continuity condition for the longitudinal components of the magnetization flux through the interface in (4.1) result in the longitudinal component \( H_{sd\parallel} \) of the magnetic field \( H_{sd} \) caused by the nonequilibrium spin polarization of spin-polarized electrons of laser-injected through the interface into an adjacent magnetic layer. The magnetic field \( H_{sd\perp} \) (independent on the spin dissipation) results in the
magnetization switching in bulk of the magnetic layer at a threshold magnitude of the laser intensity.

Thus, due to (4.1) the change of the effective magnetic field \( H_{\text{eff}} \) can result in the magnetization reorientation and switching. For the single magnetic nanolayer the effective magnetic field is caused only by the effective field \( H_{i,\text{ind}} \) related to the opto-magnetic excitations. For the tunneling magnetic junction the laser-induced effective magnetic field \( H_{i,\text{ind}} \) also includes magnetic field \( H_{sd} \) related to the laser-induced spin-polarized flux, playing the essential role in magnetization and switching processes. The last field is the sum \( H_{sd} = H_{sd,\perp} + H_{sd||} \), where the first and second terms are related to the transverse and longitudinal components of the spin flux, respectively.

The effective field \( H_{sd,\perp} \) related to the exchange s-d interaction between the lattice magnetization and the transverse component of the laser-induced spin magnetic flux damping near the magnetic interface. The effective field corresponds to the scattering of spin-polarized electrons on localized magnetic ions accompanying by the action of the torque \( T \) on the magnetic lattice. Spin magnetic momentums of the spin-polarized current and the lattice are aligned on the distance \( l \), i.e., the transverse component of the total magnetic flux is completely damped. This torque (corresponding to the continuity condition of the total magnetic flux) is determined via the spin electron polarization vector \( p_{\text{cur}} \) and magnetization vector \( m \) by the vector product

\[
T = \sigma I (m \times m \times p_{\text{cur}}) \cdot |m|, \tag{4.7}
\]

where \( \sigma \) is the constant depended on the efficiency of the scattering processes in the thin nanolayer, \( I \) is proportional to the density of the laser-induced current of spin-polarized electrons. The increase of the laser-induced current density of spin-polarized electrons to some critical value (on the order 10^7 A/cm^2) causes the large enough torque for the magnetic switching near the junction interface.

The effective field \( H_{sd||} \) is related to the longitudinal component of the total magnetic flux, consisting of laser-induced spin-polarized current and lattice magnetic components, which passes in the low-coercive layer on the spin diffusive depth (on the order 10 nm). This field generates by the exchange s-d-interaction between the non-equilibrium spin polarization and the lattice magnetization (that causes by the nonequilibrium distribution of the laser-induced electrons between spin subbands in the low-coercive magnetic layer) with the lattice magnetization. The field \( H_{sd||} \) characterizes by the direct dependence on the density of the laser-induced spin-polarized electron current. It is always parallel in the magnetization of the strongly coercive magnetic layer. Therefore, the increasing of the current density to some critical value accompanies by the increasing of \( H_{sd||} \) and magnetic switching if the magnetization directions of adjacent magnetic nanolayers are antiparallel.

4.1. The Tunnel Magnetic Memory

The nanolayer remagnetization in the tunnel magnetic junctions under the spin-polarized current and the change of last subject to the magnetization direction correspond to processes of magnetic recording and reading of an information byte, respectively. Association of these two processes is realized in the combination of the two tunnel magnetic junctions in a single magnetic nanostructure, which represents the magnetic memory element (Figure 4.1) [15].

![Figure 4.1](image.png)

Figure 4.1. Scheme of the element I and body II of magnetic memory. 1, 3, 4 are magnetic and 2 are barrier layers. Recording and reading-out are based on correlation between the spin-polarized current and a magnetization configuration of the system.

The last contains the one driven magnetic nanolayer with a weak magnetization pinning 3 sandwiched, via barrier nanolayers 2, by magnetic nanolayers with the strong magnetization pinning 1, 4. Two possible directions of the nanolayer magnetization 3 contain the information byte.

5. An Electric Field-induced Magnetization

The field-induced remagnetization in heterogeneous metal magnetic nanostructures can be realized both via laser-induced effective internal bias fields and via the electric field-induced effective surface bias fields of the Rashba spin-orbit interaction. The electric field-induced processional magnetization switching, based on the interfacial voltage-controlled magnetic anisotropy, modify the free layer perpendicular anisotropy field [15,16,17,18,19]. A bistable magnetization switching with sub-nanosecond switching time is realized by applying a unipolar voltage pulse in the FeCo/Mg/Fe magnetic tunnel junction (see [19]). Its realization in CoFeB/Mg materials system is of most technological importance for the capability of high-density integration with conventional semiconductor industry.

5.1. Spin-Orbit Influence on Surface States

For infinite crystals, the electronic band structure is related to the motion of an electron in an effective periodic potential. For finite crystals, the boundary conditions caused by the crystal surfaces, result in the occurrence of discrete surfaces states. In the framework of the density functional theory, with an effective one-particle potential, the complex values of wave vector \( \mathbf{k} \) correspond to the surface quantum states (real values of \( \mathbf{k} \) correspond to infinite crystals). The energies of the surface states lie inside the region for real \( \mathbf{k} \). Their wave functions damp in the direction of the vacuum and damp in an oscillatory way inside the crystal.
The spin-orbit interaction, that is the coupling of the orbital angular momentum and the electron spin, manifests itself via the electronic structure of solids in various ways including the magneto-crystalline anisotropy in magnets. The last makes the origin of the Rashba effects in the two-dimensional systems with broken inversion symmetry. In two-dimensional condensed matter systems (heterostructures and surface states), the combination of atomic spin-orbit coupling and asymmetry of the one-particle potential in the direction perpendicular to the two-dimensional plane causes the Rashba momentum-dependent spin splitting of energy bands. This effect can drive a wide variety of novel physical phenomena even when it is a small correction to the band structure of the two-dimensional metallic states. One is probed by angle-resolved photoelectron spectroscopy (see [20]).

The characteristic features of the Rashba spin-orbit effect, which can show up both in nonmagnetic and in magnetic heterostructures and surface states, are described in the standard model of an isotropic two-dimensional electron gas (2DEG). In this model, the effective magnetic Rashba field of the spin-orbit interaction arises in the rest frame of the moving electron and results in a Zeeman spin splitting. This field results in a unique spin topology of the electronic states both at interfaces of heterostructures and metal surfaces of at nonmagnetic and magnetic metal surfaces. The potential gradients for heterostructures and metal surfaces are determined via a band banding and an image-potential barrier, respectively.

The relativistic spin-orbit interaction and corresponding features of quantum states, the energy band structure and the spin polarization of electrons in an periodic electric field are naturally described by the Dirac equation for the four-component vector function \( \Psi \)

\[
\imath \hbar \frac{\partial \Psi}{\partial t} = H_D \Psi \tag{5.1}
\]

with the Hamiltonian

\[
H_D = \alpha \cdot \mathbf{p} + \beta mc^2 + V. \tag{5.2}
\]

Here \( \alpha = \begin{pmatrix} i & -j \end{pmatrix} \sigma \end{pmatrix} \) (\( i, j = (1, 2) \)), where \( \sigma \) is the Pauli vector-matrix, \( \mathbf{p} \) is a momentum operator; \( c \) is a light speed; the second order matrix \( \beta = \begin{pmatrix} -1 & \imath \gamma \gamma' \end{pmatrix} \),

where \( \gamma = \begin{pmatrix} 0 & 0 \\ 0 & 0 \end{pmatrix} \); \( m \) is the electron mass; \( V \) is the operator of the electron-field interaction. Stationary states are described by the wave function of the form \( \Psi = (\phi, \chi)^T \exp(\imath \mathbf{A} \cdot \mathbf{x} / \hbar) \), where \( \phi \) and \( \chi \) are two-component vector-functions.\textsuperscript{[2,3]}

Then stationary solving for Eq. (5.1) is reduced to solving the matrix equation

\[
\begin{pmatrix} E' - V & -c\sigma \mathbf{p} \\ c\sigma \mathbf{p} & E' + 2mc^2 - V \end{pmatrix} \begin{pmatrix} \phi \\ \chi \end{pmatrix} = 0, \tag{5.3}
\]

where

\[
E' = E - mc^2. \tag{5.4}
\]

Hence, in the second order approximation in \( v / c \) (\( v \) is a particle speed) it follows that

\[
\chi = \begin{pmatrix} 1 - E' - V \\ 2me^2 \end{pmatrix} \sigma \mathbf{p} \cdot \varphi. \tag{5.5}
\]

Then, eliminating the function \( \chi \) from (5.3) together with a normalization condition for \( \varphi \) give the stationary equation \( E' \varphi = H \varphi \) with the Hamiltonian

\[
H = \frac{\mathbf{p}^2}{2m} + V(r) + \frac{\hbar \sigma}{4mc} \left( \left( \nabla \times \mathbf{p} \right) \chi \right) \tag{5.6}
\]

\[
+ \frac{|E' - V(r)|^2}{2mc^2} - \frac{\hbar^2}{8mc^2} \nabla^2 V(r),
\]

where the third term represents the operator the spin-momentum interaction \( (HSO) \) for the electron in the nonuniform electric potential \( A_0 = V / e \) (\( e \) is the electron charge).

For the centrosymmetrical electric field, when \( \gamma V = \frac{e V}{r} \), this interaction take the form of the spin-orbit interaction, \( H_{SO} = \gamma_{SO}(sL) \) with the parameter

\[
\gamma_{SO} = \frac{\hbar}{2m^2c^2r} \frac{1}{dr} |V(r)|, \quad s = \sigma 2 \text{ and } L \text{ are spin and orbital moments, respectively. Taking into account the expression } m = \mu_B s \text{ for the spin magnetic moment, where } \mu_B = \frac{e \hbar}{(2mc)} \text{ is the Born magneton, the spin-orbit interaction } H_{SO} \text{ can be rewritten in the form}
\]

\[
H_{SO} = (\mu \cdot B_{SO}), \quad B_{SO} = \frac{1}{mc} |E \times p|. \tag{5.7}
\]

This magnetic field, arising in the particle’s frame, effectively couples the spin to the particle momentum, exciting the spin alignment along the field \( B_{SO} \). The spin-orbit interaction can cause a spin polarization and the spin splitting in energy.

In the two-dimensional electron system of condensed matter with the broken symmetry, the spin-orbit interaction of the one-particle potential causes the effective magnetic field \( (B_R) \) (known as the Rashba effect) inducing the electron spin polarization and the momentum dependent splitting in the electron band spectra. In general, the spin-orbit interaction parameter \( \alpha_R \) in this case depends on physical properties of the system.

In the two-dimensional free-electron condensed model, the electric potential \( A_0 \) in \( z \)-direction confines electrons to the surface. Then, from the general expression (5.2) the Hamiltonian of the system takes the form

\[
H = \frac{\hbar^2}{2m} \left( \partial_x^2 + \partial_y^2 \right) + \alpha_R (\sigma_x \partial_y - \sigma_y \partial_x), \tag{5.8}
\]

where the first term is the operator of a free electron motion in the plane \( xy \), the second term, (denoted as \( H_{RSO} \)) is the Rashba spin-orbit interaction operator in the plane \( xy \) with the parameter \( \alpha_R = \gamma V = E \). Then, two-component wave functions of the stationary wave...
function \( \Psi = (\varphi_1, \varphi_2)^T \exp(i(\omega t + k_x x + k_y y)) \) \((k_x \text{ and } k_y \text{ are components of the planar wave vector } k|| \text{ obeys the matrix equation}) \]

\[
\begin{pmatrix}
\varepsilon_k - E & -i\alpha R k_x \\
 i\alpha R k_y & \varepsilon_k - E
\end{pmatrix}
\begin{pmatrix}
\varphi_1 \\
\varphi_2
\end{pmatrix} = 0, \quad k_\pm = k_x \pm i k_y \quad (5.9)
\]

Since here, the existence of nonzero solutions assumes the zero value of the determinant of the equation matrix, then it follows the expression

\[
E_1(k) = \varepsilon_k \pm \alpha R k = \frac{\hbar^2}{2m} \left( k \pm \frac{m\alpha R}{\hbar^2} k \right)^2 - \frac{m\alpha^2}{2\hbar^2} \quad (5.10)
\]

where \( \varepsilon_k = \hbar^2 k^2 / (2m) \), describing the splitting two-dimensional electron band spectrum. The last is consisted of the two branches, caused by spin splitting under the effective Rashba field \( B_{RSO} = \frac{1}{mc} [E_z \times \sigma] \) of the spin-orbit interaction.

It is returned out, that spin polarizations of electrons with energies belonging to the different branches have opposite orientations. Indeed, the normalized spinor functions \((\varphi_1, \varphi_2)^T\) at the eigenvalues \(E_1(k)\) and \(E_2(k)\) is transformed to the spinor functions \((1, ik_x / k)^T / \sqrt{2}\) and \((ik_x / k, 1)^T / \sqrt{2}\), respectively. At \(k \parallel 0x\) these two spinor function taking the form \((1, i)^T / \sqrt{2}\) and \((i, 1)^T / \sqrt{2}\) are eigenfunctions of the spin operator \(\sigma_y\) with the eigenvalues 1 and -1, respectively. That implies that spin moments of each from both band branches for considered nonmagnetic system are opposite and lie in the plane \(xy\) as is depicted in Figure 5.1 [19].

**Figure 5.1.** The two Fermi sheets of non-magnetic metal surface corresponding two spin polarizations (denoted by arrows) of conduction electrons under the Rashba spin-orbit splitting. Here \(E\) is an energy, \(k_0\) is a splitting parameter, \(E_R\) is the spin-orbit energy. The sheets emerge from a “Dirac point” near the band bottom.

In the tight-binding model, the electrons that form the two-dimensional electron gas originate in all atomic \(s\) and \(p\) orbitals. The ingredients resulting in Rashba splitting are atomic spin-orbit interaction \(H_{SO} = \Delta_{SO} \sigma \otimes \sigma\), and an asymmetry potential in the direction perpendicular to the 2D surface \(V = E_0 z\). The main effect of the symmetry breaking potential is to open band gap \(\Delta_g\) between the isotropic \(p_z\) and \(p_x p_y\) bands. In the tight-binding approximation the hopping element from a \(p_z\) state at site \(i\) with spin \(\sigma\) to \(p_x\) or \(p_y\) state at site \(j\) with spin \(\sigma'\) is given by the expression \(t_{ij}^{x,y} = E_0 <p_z,i;\sigma | p_x,y,j;\sigma'>\). In the absence of a symmetry breaking, i.e. \(V = 0\), the hopping element vanishes due to symmetry. If \(V \neq 0\) then the hopping element is finite. The Rashba effect is the nonlinear effect, which is obtained in a second order perturbation theory with the spin-flip quantum transition between \(<p_z,i;\uparrow | <p_x,y,j;\downarrow\) states via the transitional quantum state \(<p_x,y,j;\sigma\) \). It results in the Rashba parameters \(\alpha_R = \Delta_{SO} / \Delta_g\), that in two orders of magnitude large then in the free-electron model.

### 5.1. Magnetization Via Spin-Orbit Interaction

The possibility of controlling the magnetic anisotropy of thin ferromagnetic films using an electric field \(E\) is of great interest since it can potentially lead to magnetic random access memory (MRAM) devices, which require less energy than spin-torque-transfer random access memory STT-MRAM (see [21,22]). Thin magnetic films with a perpendicular magnetic anisotropy (PMA) are important for applications. This an interfacial internal electric field might be used to engineer such a PMA is also of great interest. Experiment [23,24] has indeed shown that such a PMA might be modified by an externally applied electric field. One can be caused both by the indirect electric-induced change of magnetic anisotropy via changes of the electronic contribution to magnetic anisotropy and by the direct electric-induced changes of magnetic anisotropy via the Rashba effective interaction in combination with the exchange interaction.

The indirect electric-induced magnetization, in the terms of band theory, is related to the matrix elements of the spin-orbit interaction between empty states, large contributions to which come from regions with different \(d\)-bands (almost) cross. The strong doping dependence of these matrix elements corresponding to \(d\)-band transverse in vicinity the Fermi surface influence via the Rashba spin-orbit interaction on the magnetic anisotropy. The direct electric-induced magnetization in the framework of the Rashba spin-orbit interaction [7] and the band Stoner magnetism is characterized by a very large magnetic anisotropy arising from the internal electric fields \(E_{int}\) which exist at, e.g., ferromagnetic/metal and ferromagnetic/oxide insulator interfaces but modified by the addition of an applied electric field \(E_{ext}\). In this case the Rashba splitting there is a Rashba spin splitting of the band structure leading to a quadratic, \(E_{ext}^2\), contribution to the magnetic anisotropy, contrasting with a linear in \(E_{ext}\) doping effect.

The magnetic states in the two-dimensional model with the direct electric-induced magnetization, based on the Rashba interaction and the band Stoner magnetization, is described by Hamiltonian [19]
\[
H = \frac{p^2}{2m} - (J_0 S) \mathbf{m} \cdot \mathbf{\sigma} + \frac{\alpha_R}{\hbar} \left( \sigma_x p_y - \sigma_y p_x \right),
\]

where \( p \) is the electron momentum operator, \( \mathbf{m} = S / S \) (\( S \) is the localized spin operator), \( \mathbf{\sigma} \) is the Pauli matrices and \( \alpha_R = e \eta_{SO} \) is the Rashba parameter proportional to \( \eta_{SO} \), which characterizes the spin-orbit coupling, and the magnitude \( E \) of the electric field \( E = E z \), taken to be perpendicular to the plane of the system; \( \mathbf{m} \) is perpendicular to \( x \) and makes an angle \( \theta \) to the \( z \)-direction, as in Figure 5.2 [19].

The single \( g = S / S \), which characterizes the spin-orbit coupling, and the relative to \( \mathbf{m} \), (b) and (c) of the electric field is the same as in Eq. 5.12 but only m

\[
\mathbf{R}_R \] perpendicular to the plane, i.e., (5.14)

\[
\text{where} \ \sigma_{\alpha} = \frac{\eta_{SO} E}{\hbar} \text{ is perpendicular to the plane,} \ \mathbf{R}_R \text{ perpendicular to the plane of the system;} \ \mathbf{m} \text{ is perpendicular to } x \text{ and makes an angle } \theta \text{ to the } z \text{-direction, as in Figure 5.2 [19].}
\]

\[
\mathbf{B}_R \text{ of direction } \mathbf{k} \times \mathbf{E} \text{ lies in the } x-y \text{ plane.}
\]

For the non-magnetic two-dimensional electron system, corresponding to the surface metal nanolayer, e.g., the surface of Au, due to (5.11), the Rashba magnetic field \( \mathbf{B}_R \) results in the spin-split band energy (Figure 5.1)

\[
e_{k\sigma} = \frac{\hbar^2}{2m} - (k - \sigma k_0)^2 - E_R \quad (5.12)
\]

where the values of \( \sigma = \pm \) correspond to minority and majority electrons, respectively; the momentum shift \( k_0 = \mathbf{m} \alpha_R / \hbar^2 \), and

\[
E_R = \frac{m^2 \hbar^2}{2m^2} = \frac{1}{2} \left( -\eta_{SO} / \hbar \right)^2 \mathbf{m} E^2_R \quad (5.13)
\]

For the surface state of Au, \( E_R \approx 3.5 \text{ meV} \), exemplifying the energy scale. In the non-magnetic case (\( J_0 = 0 \)), the Rashba magnetic field as defined as \( g \mu_B \mathbf{B}_R = 2\alpha_R (-k_y x + k_x y) \), where \( g \) is the g-factor corresponding to the spin state \( |s> \) and \( \mu_B \) is the Born magneton. There are two concentric Fermi surfaces. The energy splitting \( 2\alpha_R k = \Delta (k / k_F) \), where \( \Delta \) is the value for \( k_F = (k_{F1} + k_{F1}) / 2 \), with \( k_{F1} \), the Fermi wave number for the majority/minority electrons. In the magnetic case, the magnetic order vector \( \mathbf{m} = \cos \theta z + \sin \theta y \). The total Rashba field \( \mathbf{B}_{Rsd} \) defining the axis of quantization,

\[
g \mu_B \mathbf{B}_{Rsd} = 2 \left[ (J_0 S + \alpha_R k_x \sin \theta) + \alpha_R (-k_y x + k_x \cos \theta (m \times x)) \right] \quad (5.14)
\]

due to (5.11), results in the non-symmetrical spin-split energy band

\[
e_{k\sigma} = \frac{\hbar^2}{2m} \left[ (k_x - \sigma k_0 \sin \theta)^2 + k_y^2 \right] - E_R \sin \theta \quad \sigma = 1, \ -1
\]

\[\sigma = \left( J_0 S \right)^2 + \alpha_R^2 \left( k_x^2 \cos^2 \theta + k_y^2 \right)^{1/2}.\]

The direction of the momentum shift depends upon \( \sigma \), i.e., the majority/minority character of the band. It is assumed, that \( g \mu_B B_R < J_0 S \), i.e. the Rashba magnetic field \( B_R \) is smaller than the exchange splitting. The second order in \( g \mu_B B_R \) and \( g \mu_B B_{Rsd} = 2JS \), where \( JS = \left[ (J_0 S)^2 + \alpha_R^2 (k_x^2 \cos^2 \theta + k_y^2) \right]^{-1/2} \) / \( J_0 S \) and where \( m' \) differs from \( m \) by a small angle \( \delta = \tan^{-1} (\alpha_R k_x^2 \cos^2 \theta + k_y^2) / (J_0 S) \).

The contributions to the magnetic anisotropy are highlighted by contrasting the perpendicular and parallel orientations of magnetic order vector \( m \) to the plane. With \( m \) perpendicular to the plane, i.e., \( m \parallel z \) (\( \theta = 0 \)), the exchange and Rashba fields are orthogonal and hence the net energy for a single electron (5.15) is

\[
e_{k\sigma} = \frac{\hbar^2}{2m} k^2 - \sigma \left[ (J_0 S)^2 + (\alpha_R k^2) \right]^{1/2}. \quad (5.16)
\]

The axis of quantization is tilted by \( \delta(k) = \tan^{-1}(\alpha_R k_0) / (J_0 S) \) away from the z -axis (Figure 5.3(a)).

On the other hand, the \( x \)-component of \( \mathbf{B}_R \), which is perpendicular to \( J_0 S \), gives rise to a correction to the effective exchange field. The direction of the moment tilts away from the y-axis in the direction perpendicular to the wave vector by \( \delta(k_\gamma) = \tan^{-1}(ak_\gamma) / (J_0 S) \). The single particle energy is described by the expression,

\[
e_{k\sigma} = \frac{\hbar^2}{2m} \left[ (k_x - \sigma k_0 \sin \theta)^2 + k_y^2 \right] - E_R \quad (5.17)
\]

\[-\sigma \left( J_0 S^2 + \alpha_R^2 k^2 \right)^{1/2} \]

\[\text{where the shift } k_0 \text{ is the same as in Eq. 5.12 but only along the } x\text{-axis. The effective exchange field in Eq. 5.17 is smaller than that in Eq. 5.15 due to the absence of a } k_\gamma^2 \text{ term. This indicates that the overall DM contribution favours a perpendicular } \mathbf{m} \text{ while the PD term favours an in-plane } \mathbf{m} \text{. This exchange field changes sign for the majority/minority spins, i.e., with } \sigma \text{. Assuming}
\]
\((J_0 S)^2 > (\alpha_R k_y)^2\) and retaining the \(\theta\)-dependent terms up to the order of \(E^2\) in (5.14), we obtained

\[ E_{an} = E_R \left[ 1 - \frac{2T}{J_0 S} \right] \cos^2 \theta \]  
(5.18)

for the magnetic anisotropy energy, with

\[ T = \frac{\hbar^2}{2m} \left\langle k_x^2 \right\rangle \left\langle \sigma \right\rangle \left\{ \sigma \right\} \left\{ \sigma \right\} \langle \downarrow \rangle \],

(5.19)

where \(\left\{ \right\}\) denotes an average over the Fermi surface. The Rashba spin-orbit interaction produces an uniaxial anisotropy energy, which, as in the Dzyaloshinskii-Moriya theory, comprises a direct second order in \(E\) easy plane pseudo-dipolar interaction and indirect contribution proportional to \(E^2/J_0 S\). Clearly an \(E^2\) dependent PMA results, if \(T > J_0 S/2\), which is the case for real 3d ferromagnets.

The resulting anisotropy energy can be very large. The value of the scaling pre-factor \(E_R\) in Eq. (5.18) for the surface state of Au is \(\sim 3.5\) meV or about 35 T in magnetic field units and very much larger than the typical \(\sim 1\) T demagnetising field. If the Au film is polarized by contact with an ultra-thin ferromagnet the second factor, \(2T/J_0 S\), in Eq. (5.17) for the field inside a Au surface layer can be quite large \(\sim 5\) leading to a PMA and indeed ultra-thin Fe on Au does have such a PMA [25].

Schematically shown in Figure 5.4(a) [19] is the potential seen by electrons in a freestanding ultra-thin ferromagnetic film.

Figure 5.4. (a) There is an electric field \(E\) in the surface region of a ferromagnet, however for a given wave vector \(k\), the Rashba field \(B_R\), proportional to \([k \times E]\), has an opposite sign at the two surfaces and the average field is zero. (b) With a finite external field this symmetry is broken and there is a net Rashba field acting upon the electrons.

The electric field is increased at one surface and is decreased at the other doubling the net effect. In contrast, for this same symmetric situation, the surface charges are opposite and doping effects must cancel. Clearly the intrinsic Rashba field \(B_R\) is modified when the material adjacent to the ferromagnets (F), say Fe, are different. In many experiments an insulator I, often MgO, lies to one side and a normal metal (N), and e.g., Au, Pt or Pd, completes a tri-layer system.

Thus, the Rashba effective magnetic field due to the internal electric field in the surface region of nanolayered ferromagnetics can make an important contribution to the perpendicular magnetic anisotropy. The Rashba spin-orbit interaction in combination with the exchange interaction between spin-polarized electrons and the localized lattice magnetization allow the electric field magnetization control in in magnetic nanostructures including multilayered magnetic nanostructures.

References


