American Journal of Nano Research and Applications
2017; 5(5): 69-80
http://www.sciencepublishinggroup.com/j/nano
doi: 10.11648/j.nano.20170505.12
ISSN: 2575-3754 (Print); ISSN: 2575-3738 (Online)

Review Article

Spin-Dependent Currents in Magnet/Normal Metal Based Magnetic Nanostructures

Andrii Korostil*, Mykola Krupa

Department of Magnetic Materials and Nanocrystalline Structures, Institute of Magnetism NASU, Kyiv, Ukraine

Email address:
amand@rambler.ru (A. Korostil)
*Corresponding author

To cite this article:
Andrii Korostil, Mykola Krupa. Spin-Dependent Currents in Magnet/Normal Metal Based Magnetic Nanostructures. *American Journal of Nano Research and Applications.* Vol. 5, No. 5, 2017, pp. 69-80. doi: 10.11648/j.nano.20170505.12

Received: October 10, 2017; **Accepted**: October 24, 2017; **Published**: November 27, 2017

Abstract: The spin transport through and near interfaces have been studied in magnet/normal metal based multilayer magnetic nanostructures in magneto-static and magneto-dynamic cases. Its features and accompanying effects, such as the magnetoresistance or the magnetic precession induced spin pumping and spin accumulation in adjacent normal metal are determined by the spin-dependent scattering on the interface. These effects are governed by the entire spin-coherent region that is limited in size by spin-flip relaxation processes and can be controlled by the spin-polarized current of different origin including the spin Hall effect. Conditions of realization of the mentioned spin currents in the multilayer magnetic nanostructures are studied.

Keywords: Spin Currents, Magnetic Nanostructures, Spin-Dependent Scattering, Magnetic Dynamics, Spin Pumping

1. Introduction

Coupling between spin currents and localized magnetic moments in magnet (M)/normal metal (NM) based multilayer magnetic nanostructures constitutes the basis of the mutual control between electric current and static or dynamic magnetic states. Herewith, magnetic layers include magnets with the exchange interactions both ferromagnetic (F) and antiferromagnetic (AF) types (e.g., ferrimagnetics compounds like YIG, Gd_3 Fe_5O_{12} and AFs Fe_3O_4, $NiFe_2O_4$, NiO [1, 2]), normal metalsare nonmagnetic, usually, heavy metals with strong spin-orbit coupling (e.g., Pt, Ta, W). The mentioned interconnection in these magnetic nanostructures occur via the interface scattering of the spin-polarized current and its *s-d* exchange interaction with static or dynamic magnetic states [3-6]. The impact of the spin current on the magnetic states is manifested through the spin-transfer torque and the impact of the localized magnetic momentum on the spin current is manifested via the spin dependent interface scattering accompanying by magnetoresistance effect. The spin polarization can be induced by effective bias fields of different origin including fields caused by an exchange interaction and the strong spin-orbit coupling. The entire spin-coherent region is limited in size by spin-flip relaxation processes.

In the case of static magnetic states, the mutual influence of the spin current and magnetic ordering can be manifested as the magnetiresistance effect of the dependence of the spin current on the magnetization orientation in the magnetic layer and vice versa, the dependence of the latter the spin current [5, 6]. Such effects can constitute the base for magnetic writing techniques in non-volatile memory technologies such as MRAM [7] and racetrack memories [8]. They also include the giant magneto-resistance (GMR) effect in metallic magnetic multilayers, which has commercially utilized in high-end magnetic recording media [9]. Obtaining the mentioned multilayer magnetic nanostructures with properties of electric-controlled magnetic switching and the magnetic-controlled spin current involves the description of features of the spin transport in magnetic heterogeneous nanostructures allowing for the compatibility conditions at the interfaces [4, 10]. This is usually solved within the Landauer-Büttiker formalism [11] and more rigorously, using the non-equilibrium Keldysh-Green functions [12, 13].

In the case of the dynamic magnetic states, their

interconnection with the control spin current is affected by the magnetic precession-induced spin pumping and the spin accumulation in the normal metal layer at the interface [3, 10]. The action of the spin currents on the magnetic dynamics via the spin-transfer torque and the reciprocal process of spin pumping result in the effect of controlled magnetic auto-oscillations [14]. The magnetic dynamic damping is related to the spin pumping effect at the M|NM interface that can be compensated by the spin-transfer torque from the spin current of the converted input current. This spin transfer is governed by the reflection and transmission matrices of the system, analogous to the scattering theory of transport and interlayer exchange coupling. Due to interfacial processes, M|N coupling becomes important in the limit of ultrathing (≤ 10 nm) magnetic films and can lead to a sizable enhancement of the damping constant.

The above-mentioned coupling effects at interfaces can occur in the magnetic nanostructures with both ferromagnetic (F) and antiferromagnetic (AF) exchange interactions, which are realized in fero- ferri- and antiferromagnetic materials. Normal metal layers are medium for the spin currents, which can be converted from the control charge current by the spin-orbit interaction, especially, the spin Hall and the spin-orbit Rashba [15, 16] effects.

The paper is organized as follows. In Sec. 1 the spin-dependent transport in the F/N based magnetic nanostructure is studied for the static magnetization. In the modified Stoner model with potential barrier dependent on the physical parameters including the magnetization directions, the chemical potentials of the layers, and the contact conductances, the parametrically dependent scattering of spin-polarized current is investigated. The mentioned parameters are determined by the spin-polarized kinetic equations in the framework of the Keldysh Green function approach. It is considered both in single and composite F/N based magnetic nanostructures. In Sec. ll features of the interconnection between magnetization dynamics and the spin currents are studied in the F/N based nanostructures. The process of the magnetization precession-induced pumping spin current in the nonmagnetic layers is considered as the result of the parametric time dependence of the interfacial scattering with the precession as the parameter. It is shown that the spin pumping slows down the precession corresponding to an enhanced Gilbert damping constant in the Landau-Lifshitz-Gilbert model. The spin current related to the spin pumping, which flows back into the ferromagnetic layers and driven by the accumulated spins in the normal metal layers is also discussed.

2. Spin Transport in the Case of Static Magnetic Field

2.1. Features of Spin-Dependent Electric Current in the F/N Belayers

Characteristic features of the spin-dependent transport and the interfacial scattering in multilayer magnetic nanostructures based on F/N bilayers are manifested the F/N bilayer (Figure 1). These features are related to the conditions under which long-range spin effects are observable in normal metals. Spins injected into a normal metal layer relax due to unavoidable spin-flip processes.

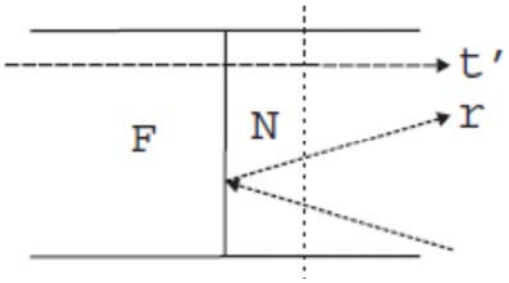

Figure 1. *A contact between a ferromagnetic (F) and a normal (N) metal layers. At the normal metal side, the current is denoted as the dotted line. The transmission coefficient from the ferromagnet to the normal metal is t and the reflection matrix from the normal metal to the normal metal is r.*

Naturally, the dwell time on the layer must be shorter than the spin-flip relaxation time in order to observe nonlocality in the electron transport. For a simple ferromagnet (F) normal metal (N) double heterostructure (F/N/F) with antiparallel magnetizations the condition can be quantified following [17]. The spin-current into the normal metal layer is roughly proportional to the particle current, $e(ds/dt)_{tr} \sim I = V/R$, where s is the number of excess spins on the normal metal layer, V is the voltage difference between the two reservoirs coupled to the normal metal layer, and R is the F|N contact resistance. When the layer is smaller than the spin-diffusion length, the spin-relaxation rate is $e(ds/dt)_{rel} = -s/\tau_{sf}$, where τ_{sf} denotes the spin-relaxation time on the layer. (Otherwise, this simple approach breaks down since the spatial dependence of the spin-distribution in the normal metal should be taken into account [18]). The number of spins on the normal metal layer is equivalent to a non-equilibrium chemical potential difference $\Delta\mu = s\delta$ in terms of the energy level spacing δ (the inverse density of states) (more generally the relation between $\Delta\mu$ and s is determined by the spin-susceptibility). The spin-accumulation on the normal metal layer significantly affects the transport properties when the non-equilibrium chemical potential difference is of the same order of magnitude or larger than the applied source-drain voltage, $\Delta\mu > eV$ or $\delta\tau_{sf}/\hbar > R/R_K$, where $R_K = e^2/h$ is the quantum resistance. Thus, spin-accumulation is only relevant for sufficiently small normal metal layers and/or sufficiently long spin-accumulation times and/or good contact conductances.

The spin-dependent current in the model F/N bilayer (Figure 1) is expressed via the 2×2 distribution matrix matrix $f(\varepsilon)$ in spin-space at a given energy ε in the layer. The external reservoirs are assumed to be in local equilibrium so that the distribution matrix is diagonal in spin-space and attains its local equilibrium value $f = 1f(\varepsilon, \mu_\alpha)$, where 1 is the unit matrix, $f(\varepsilon, \mu_\alpha)$ is the Fermi- Dirac distribution

function and μ_α is the local chemical potential in reservoir α. The direction of the magnetization of the ferromagnetic layers is denoted by the unit vector $\boldsymbol{m}_\alpha$.

The 2×2 non-equilibrium distribution matrices in the layers in the stationary state are uniquely determined by current conservation

$$\sum_\alpha I_{\alpha\beta} = \left(\frac{\partial f_\beta}{\partial t}\right)_{rel} \tag{1}$$

where $I_{\alpha\beta}$ denotes the 2 × 2 current in spin-space from layer (or reservoir) α to layer (or reservoir) β and the term on the right hand side describes spin-relaxation in the normal layer. The right hand side of the equation (1) can be set to zero when the spin-current in the layer is conserved, i.e. when an electron spends much less time on the layer than the spin-flip relaxation time τ_{sf}. If the size of the layer in the transport direction is smaller than the spin flip diffusion length $l_{sf} = \sqrt{D\tau_{sf}}$, where D is the diffusion coefficient then the spin-relaxation in the layer can be introduced as $\left(\partial f^N / \partial t\right)_{rel} = \left(1\mathrm{Tr}(f^N)/2 - f^N\right)/\tau_{sf}$.

2.2. Passing the Electric Current Through the F|N Contact

The relation between the current in the F/N bilayer through the F|N interface and electron distribution functions in F and N can be described by the non-equilibrium Keldysh-Green functions and their equations representing the quantum-kinetic equations for the one-particle propagator. The corresponding Hamiltonian has the form

$$H = h(1) + w(1,2) \tag{2}$$

where

$$h(1) = -\frac{1}{2m}\left[\left(-i\hbar\nabla + A(1)\right)^2 + \varphi(1)\right] + V^p(1) + V^s(1)$$

is the one-particle operator describing electrons in electromagnetic field, $V^p(1)$ and $V^s(1)$ are spin-independent and spin-dependent parts of the one-electron potential, $w(1,2)$ is the operator of the two-particle interaction. Here 1 and 2 denote coordinates $\boldsymbol{r}_1$ and $\boldsymbol{r}_2$. The spin-dependent potential $V^s(1) = (\sigma \cdot \boldsymbol{m})W^s(1)$ existences only in the magnetic metal and vanishes in the normal metal.

The physical properties of the system are described by the one-particle non-equilibrium Green function, i.e., the expectation value of the time contour-ordered product of creation and annihilation field operators $\psi^\dagger(1)$ and $\psi(1)$, respectively, with $1 = (r_1, s_1; z_1)$ (s_1 is the spin variable). The variable z_i changes along the Keldysh contour, from the initial real time t_0 (before which the system is in an equilibrium state) to t, that correspond to the forward time-evolution of the state, and then it change from t to t_0, that correspond to backward time-evolution of the system. The change of z_i on the imaginary time interval $[t_0, t_0 - i\beta]$ (where β is the inverse temperature) corresponds to the equilibrium state. The conjunction of the mentioned time intervals forms the so-called Keldysh contour C consisting of forward and backward real-time branches and the thermal imaginary track. This non-equilibrium Green function can be presented as [12, 13]

$$G(1,2) = -i\frac{\mathrm{Tr}\left[T_C e^{-i\int_C H(z)\,\mathrm{d}z}\psi(1)\psi^\dagger(2)\right]}{\mathrm{Tr}\left[T_C e^{-i\int_C H_0(z)\,\mathrm{d}z}\psi(1)\right]},$$

$$= -i\left\langle \psi_H(1)\psi_H^\dagger(2)\right\rangle \tag{3}$$

where T_C is contour time-ordering operator on the Keldysh contour C with $H_0(z) = H(z)|_{z\in[t_0, t_0 - i\beta]}$, the subscript H in the right hand side denotes the Heisenberg representation. Due to (3) the Green function can be presented as

$$G(1,2) = \theta(z_1, z_2)G^>(1,2) + \theta(z_2, z_1)G^<(1,2) \tag{4}$$

where $\theta(z_1, z_2)$ is the contour step function equal 1 or 0 versus z_1 or z_2 is later on the contour C and the greater ($G^>$) and lesser ($G^<$) Green functions are determined by the relations

$$G^>(1,2) = -i\left\langle \psi_H(1)\psi_H^\dagger(2)\right\rangle,\ G^<(1,2) = i\left\langle \psi_H^\dagger(2)\psi_H(1)\right\rangle \tag{5}$$

where an extra – sign is introduced due to the interchange of the Fermionic operators by the contour ordering operator T_C. The Green function provides a direct access to observable physical quantities of the system. For example, the equal-time limit gives directly the particle spin density at the space-time point 1

$$\left\langle n(1)\right\rangle = \left\langle \psi_H^\dagger(1)\psi_H(1)\right\rangle = -G^<(1,1^+) \tag{6}$$

(a superscript "+" means infinitesimal). The spin current density is determined as

$$I(1) = -\left\{\partial_t G^<(t,t') + \partial_{t'}G^<(t,t')\right\}_{t'=t} \tag{7}$$

The Green functions are described by the equation

$$\left\{i\partial_{z_1} - h(1)\right\}G(1,2) = \delta(1,2) + \int d\,3w(1,3)G(1,3;2,3^+) \tag{8}$$

following of the Schrödinger equation for wave functions of the system. Here, the two-particle Green function

$$G(1,3;2,3^+) = \left\langle \psi_H^\dagger(2)\psi_H(1)\psi_H(3)\psi_H^\dagger(3^+)\psi_H^\dagger(2)\right\rangle \tag{9}$$

expresses via the functional derivative with respect to the variation in the infinitesimal external potential ν of the one-particle Green function by the relation [12]

$$G(1,3;2,3^+) = i\left\langle n(3)\right\rangle G(1,2) + \frac{\delta G(1,2)}{\delta v(3)} \tag{10}$$

Consequently, the equation (8) can be represented as self-contained functional derivative equation with the matrix representation

$$\left(L^0 + L^1\right) G = 1 \tag{11}$$

where the matrix L^0 is determined by matrix elements

$$L^0(1,2) = \left(i\partial_z - h^*(1) + \int d3\, w(1,3)\left\langle n(3)\right\rangle\right)\delta_{1,2} \tag{12}$$

not containing the functional derivative. The matrix L^1 is determined by matrix elements

$$L^1(1,2) = i\delta_{1,2}\int d3\, w(1,3)\frac{\delta}{\delta v(3)} \tag{13}$$

which proportional to the functional derivative.

The expression for the functional derivative

$$\frac{\delta}{\delta v(3)} G = \frac{\delta}{\delta v(3)} G\left[G^{-1}\right] G \tag{14}$$

(the bracket separates the expression experiencing the variation differentiation) allows to represent the operation of the differential matrix L^1 on G via the self-energy matrix Σ not containing functional derivatives:

$$L^1 G = \Sigma G, \quad \Sigma = L^1 G\left[G^{-1}\right] \tag{15}$$

Finally, the matrix equation (10) reduces to the system

$$\left(L^0 + \Sigma\right) G = 1 \tag{16}$$

$$\Sigma = -L^1 G\left[L^0 + \Sigma\right] \tag{17}$$

The first-order approximation with respect to the interaction w determines the Hartree-Fock self-energy matrix $\Sigma^{\mathrm{HF}} = -L^1 G\left[L^0\right]$ with matrix elements

$$\Sigma^{\mathrm{HF}}(1,2) = -\delta(1,2)\int d3\, w(1,3)\left\langle n(3)\right\rangle + iw(1,2)G(1,2) \tag{18}$$

Here the first term describes the classical Hartree potential at 1 produced by the charge density throughout the space and the second term is the space-nonlocal exchange potential originating from the Pauli exclusion principle and antisymmetry of the wave functions. Due to (16) the self-energy matrix in second-order approximation is determined by the equation

$$\Sigma = -L^1 G\left[L^0\right] + L^1 G\left[L^1 G\left[L^0\right]\right] \tag{19}$$

with matrix elements of the form

$$\Sigma^2 = \Sigma^{\mathrm{HF}}(1,2) + G(1,2)\int d3d4\, w(1,3)w(2,4)G(4,3)G(3,4^+)$$
$$- \int d3d4\, G(1,3)w(1,4)G(3,4)G(4,2)w(3,2) \tag{20}$$

Here the second and third terms describe the correlation and scattering effects.

Equation Chapter 1 Section 1 In the case under considered of negligible small influence of correlation and scattering effects in the stationary situation, the Green function in the energy representation for the bilayer Fe/N is decomposed into quasi-one-dimensional modes as

$$G_{ss'}(1,1') = \sum_{nm,\alpha\beta} G^{\alpha\beta}_{nsms'}\chi^n_s(\rho,x)\chi^{m*}_{s'}(\rho',x')e^{i\alpha k^n_s x - i\beta k^m_{s'}x'} \tag{21}$$

Here the indices $\alpha,\beta = (+,-)$, where the signs "+" and "−" denote right-going (+) and left-going (−) modes, respectively, $\chi^n_s(\rho,x)$ is the transverse wave function and k^n_s denotes the longitudinal wave-vector for an electron in transverse mode n with spin s. Then, based on the definition of the current through the Green function and taking into account the spatial independent the transverse wave function $\chi^n_s(\rho,x)$ can be obtained the expression

$$I_{ss'}(x) = ie\sum_{n\alpha\beta}\left(\alpha v^n_s - \beta v^m_{s'}\right)G^{\alpha\beta}_{nsms'}(x,x)\int d\rho\chi^n_s(\rho,x)\chi^m_{s'}(\rho,x) \tag{22}$$

for the spin current, where $v^n_s = \hbar k^n_s / m$ is the longitudinal velocity for an electron in transverse mode n with spin s. In a normal metal, the transverse states and the longitudinal momentum are spin-independent and the spin current simplifies to

$$I_{ss'}(x) = 2ie\sum_{n\alpha}\alpha v^n_s G^{\alpha\alpha}_{nsms'}(x,x) \tag{23}$$

which is used in the calculation of the spin current on the normal side of the contact.

Using the representation

$$G^{\alpha\beta}_{nsms'}(x,x') = -i\left(\frac{g^{\alpha\beta}_{nsms'}(x,x')}{\sqrt{v^n_s v^m_{s'}}} + 1\delta_{ss'}\frac{\alpha\delta_{\alpha,\beta}\,\mathrm{sign}(x-x')}{v^n_s}\right) \tag{24}$$

where the latter term does not contribute to the current on the normal side of the interface F|N, the expression

$$I_{ss'}(x) = 2e\sum_{n\alpha}\alpha v^n_s g^{\alpha\alpha}_{nsms'}(x,x) \tag{25}$$

can be obtained for the spin current on the normal metal side.

The complete description of the spin current through the F|N interface involves taking into account the connection between waves propagating to the right (left) on the right hand side of the contact ψ^+_R (ψ^-_R) and waves propagating to the right (left) on the left hand side of the contact ψ^+_L (ψ^-_L). This is described by the transfer matrix M obeying the relation

$$\begin{pmatrix} \psi_R^+ \\ \psi_R^- \end{pmatrix} = M \begin{pmatrix} \psi_L^+ \\ \psi_L^- \end{pmatrix} \tag{26}$$

which in terms of the transmission and reflection coefficients takes the form

$$M = \begin{pmatrix} t - r'(t')^{-1} r & r'(t')^{-1} \\ -(t')^{-1} r & (t')^{-1} \end{pmatrix} \tag{27}$$

Here the transmission and reflection coefficients enter in definition of the scattering matrix

$$S = \begin{pmatrix} r & t' \\ t & r' \end{pmatrix} \tag{28}$$

where $r_{nm}^{s\sigma}$ is the reflection matrix for incoming states from the left in mode m and spin σ to mode n with spin s, $t_{nm}^{s\sigma}$ is the transmission matrix for incoming states from the left transmitted to outgoing states to the right. In addition, r' is the reflection matrix for incoming states from the right reflected to the right, and t' is the transmission matrix for incoming states from the right transmitted to the left. The Green function to the left ($x = x_2$) of the interface is expressed via the Green function to the right ($x = x_1$) of the interface by the expression

$$g_{nsms'}^{\sigma\sigma'}(x = x_2, x') = \sum_{ls'',\sigma''} M_{nsls''}^{\sigma\sigma''} g_{ls''ms'}^{\sigma''\sigma'}(x = x_1, x') \tag{29}$$

In the matrix form $g_2 = M g_1 M^\dagger$, ($g_{2(1)} = g(x = x_{2(1)}$, $x' = x_{2(1)})$).

In the approximation of isotropic quasi-classical Green functions in nanolayers 1 and 2, $(G_1)_{nsms'}^{\alpha\beta} = \delta_{n,m}\delta^{\alpha\beta}(G_1)_{ss'}^{\alpha\beta}$. Considering the representation of the Green function in terms of the retarded G^R, advanced G^A and Keldysh G^K Green functions

$$G = \begin{pmatrix} G^R & G^K \\ 0 & G^A \end{pmatrix}$$

it can be obtained that $G_R = -G_A = -1$ and $G_{K,1(2)} = h_{1(2)}1$, where the two-dimensional matrix h is related to the non-equilibrium distribution functions $h_{1(2)} = 2(f(\varepsilon)_{1(2)} - 1)$. Herewith

$$g_{R,1}^{-+} = 2r,\ g_{R,2}^{+-} = 2r' \tag{30}$$

and

$$g_{K,1}^{--} = t' h_2 t'^\dagger + h_1 r^\dagger,\ g_{K,2}^{--} = t h_1 t^\dagger + r' h_2 r'^\dagger \tag{31}$$

Inserting the expressions (31) into (25) results in the expression

$$I = \frac{e}{h}\left\{\sum_{mm}\left[t'^{nm} f^F (t'^{mn})^+ - \left(M f^F - \tau^{nm} f^N (\tau^{mn})^+\right)\right]\right\} \tag{32}$$

which describes the current through interface on its normal metal side. Here $r_{ss'}^{nm}$ is the reflection coefficient for electrons from transverse mode m with spin s' incoming from the normal metal side reflected to transverse mode n with spin s on the normal metal side, and $t'^{nm}_{ss'}$ is the transmission coefficient for electrons from transverse mode m with spin s' incoming from the ferromagnet transmitted to transverse mode n with spin s on the normal metal side. (Note that the hermitian conjugate in (32) operates in the spin-space and the space spanned by the transverse modes, e.g. $(r^{mn})_{ss'}^+ = (r^{nm})_{s's}^*$).

2.3. Parametric Spin Dependence of Electric Current Through Contact

The relation (32) between the current and the distribution functions has a simple form after transforming the spin-quantization axis. Disregarding spin-flip processes in the contacts, the reflection matrix for an incoming electron from the normal metal transforms as

$$r^{nm} = \sum_s u^s r_s^{nm} \tag{33}$$

where $r_{\uparrow(\downarrow)}^{nm}$ are the spin-dependent reflection coefficients in the basis where the spin-quantization axis is parallel to the magnetization in the ferromagnet, the spin-projection matrices are

$$u^{\uparrow(\downarrow)} = (1 \pm \sigma \cdot \boldsymbol{m})/2 \tag{34}$$

and σ is a vector of Pauli matrices. Similarly for the transmission matrix $t_{\uparrow(\downarrow)}^{nm}$

$$t'^{nm}(t'^{mn})^+ = \sum_s u^s \mid t_s'^{nm} \mid^2, \tag{35}$$

where $t'^{nm}_{\uparrow(\downarrow)}$ are the spin-dependent transmission coefficients in the basis where the spin-quantization axis is parallel to the magnetization in the ferromagnet. From the unitarity of the scattering matrix, it is follows that the general form of the relation (32) reads

$$eI = G^\uparrow u^\uparrow \left(f^F - f^N\right) u^\uparrow + G^\downarrow u^\downarrow \left(f^F - f^N\right) u^\downarrow$$

$$-G^{\uparrow\downarrow} u^\uparrow f^N u^\downarrow - \left(G^{\uparrow\downarrow}\right)^* u^\downarrow f^N u^\uparrow \tag{36}$$

where it is introduced the spin-dependent conductances $G^\uparrow$ and $G^\downarrow$

$$G^{\uparrow(\downarrow)} = \frac{e^2}{h}\left[M - \sum_{nm} \mid r_{\uparrow(\downarrow)}^{nm} \mid^2\right] = \frac{e^2}{h}\sum_{nm} \mid t_{\uparrow(\downarrow)}^{nm} \mid^2 \tag{37}$$

and the mixing conductance

$$G^{\uparrow\downarrow} = \frac{e^2}{h}\left[M - \sum_{nm} r_{\uparrow}^{nm}\left(r_{\downarrow}^{nm}\right)^*\right] \tag{38}$$

The precession of spins leads to an effective relaxation of spins non-collinear to the local magnetization in ferromagnets and consequently the distribution function is limited to the form $f^F = 1f_0^F + \sigma\cdot\boldsymbol{m}\, f_s^F$. Such a restriction does not appear in the normal metal layer and f^N can be any hermitian 2 × 2 matrix.

The relation between the current through a contact and the distributions in the ferromagnetic layer and the normal metal layer are determined by four parameters, the two real spin-dependent conductances ($G^{\uparrow}, G^{\downarrow}$) and the real and imaginary parts of the mixing conductance $G^{\uparrow\downarrow}$. These contact-specific parameters can be obtained by microscopic theory or from experiments. The spin-conductances $G^{\uparrow}$ and $G^{\downarrow}$ describe spin-transport for a long time [19]. The mixing conductance is relevant for transport between non-collinear ferromagnets. Note that although the mixing conductance is a complex number the 2 × 2 current in spin-space is hermitian and consequently the current and the spin-current in any direction given by (36) are real numbers. Due to the definitions of the spin-dependent conductances (37) and the 'mixing' conductance (38)

$$2\,\mathrm{Re}\,G^{\uparrow\downarrow} = G^{\uparrow} + G^{\downarrow} + \frac{e^2}{h}\sum_{nm}\left|\, r_{\uparrow}^{nm} - r_{\downarrow}^{nm}\,\right|^2 \tag{39}$$

and consequently, the conductances should satisfy the relation $2\,\mathrm{Re}\,G^{\uparrow\downarrow} \geq G^{\uparrow} + G^{\downarrow}$.

In terms of a scalar particle and a vector spin-contribution $\boldsymbol{I} = (1\boldsymbol{I}_0 \pm \sigma\cdot\boldsymbol{I}_s)/2,\quad \boldsymbol{f}^{N(F)} = 1\boldsymbol{f}_0^{N(F)} \pm \sigma\cdot\boldsymbol{s}_s(\boldsymbol{m})\boldsymbol{f}_s^{N(F)}$ and the particle current is described by the expression

$$I_0 = \left(G^{\uparrow} + G^{\downarrow}\right)\left(f_0^F + f_0^N\right) + \left(G^{\uparrow} - G^{\downarrow}\right)\left(f_s^F \boldsymbol{\cdot m}\cdot\boldsymbol{s}\, f_s^N\right) \tag{40}$$

The familiar expressions for collinear transport are recovered when $\boldsymbol{m}\cdot\boldsymbol{s} = \pm 1$. The spin-current is

$$\begin{aligned}\boldsymbol{I}_s = \boldsymbol{m}\Big[&\left(G^{\uparrow} - G^{\downarrow}\right)\left(f_0^F - f_0^N\right) + \left(G^{\uparrow} + G^{\downarrow}\right) f_s^F \\ &+ \left(2\,\mathrm{Re}\,G^{\uparrow\downarrow} - G^{\uparrow} - G^{\downarrow}\right)\boldsymbol{m}\cdot\boldsymbol{s}\, f_s^N\Big] \\ &-2\boldsymbol{s}\,\mathrm{Re}\,G^{\uparrow\downarrow} f_s^N + (\boldsymbol{m}\cdot\boldsymbol{s})2\,\mathrm{Im}\,G^{\uparrow\downarrow} f_s^N\end{aligned} \tag{41}$$

The first three terms point in the direction of the magnetization of the ferromagnet $\boldsymbol{m}$, the fourth term is in the direction of the non-equilibrium spin-distribution $\boldsymbol{s}$, and the last term is perpendicular to both $\boldsymbol{s}$ and $\boldsymbol{m}$. The last contribution solely depends on the imaginary part of the mixing conductance. This term can by interpreted by considering how the direction of the spin on the normal metal layer s would change in time keeping, all other parameters constant. The cross product creates a precession of $\boldsymbol{s}$ around the magnetization direction $\boldsymbol{m}$ of the ferromagnet similar to a classical torque while keeping the magnitude of the spin-accumulation constant. In contrast, the first four terms represent diffusion-like processes, which decrease the magnitude of the spin-accumulation. Due to (40) the non-equilibrium spin-distribution f_s^N propagates easier into a configuration parallel to $\boldsymbol{s}$ than parallel to $\boldsymbol{m}$, since these processes are governed by positive diffusion-like constants $2\,\mathrm{Re}\,G^{\uparrow\downarrow}$ and $2\,\mathrm{Re}\,G^{\uparrow\downarrow} - G^{\uparrow} - G^{\downarrow}$, respectively.

2.4. Spin-Dependence Currents for Different Types of Contacts

The four conductance parameters $G^{\uparrow}$, $G^{\downarrow}$, $\mathrm{Re}\,G^{\uparrow\downarrow}$ and $\mathrm{Im}\,G^{\uparrow\downarrow}$ depend on the microscopic feature of interfaces at given the crossection A, the length of the normal and ferromagnetic parts of the contact L^N and L^F, respectively, the conductivity on the normal side σ^N and the spin-dependent conductance of the ferromagnetic part σ^{F_s}. They are different for a diffusive, a ballistic, and a tunnel contact.

In the case a diffusive contact between a normal metal and a ferromagnet, the conductances of the normal and ferromagnetic parts are $G_D^N = A\sigma^N / L^N$ and $G_D^F = A\sigma^F / L^F$, respectively. The spin-dependent conductances of the whole contact are obtained as the diffusive ferromagnetic and normal metal regions:

$$G_D^{\uparrow} = \frac{G_D^{F\uparrow} G_D^N}{G_D^{F\uparrow} + G_D^N},\; G_D^{\downarrow} = \frac{G_D^{F\downarrow} G_D^N}{G_D^{F\downarrow} + G_D^N} \tag{42}$$

These spin-dependent conductances ($G_D^{\uparrow}$ and $G_D^{\downarrow}$) fully describe collinear transport (in the absence of spin-flip scattering). For non-collinear magnetizations the mixing conductance, which is also needed, can be derived from the scattering matrix. The latter follows from the diffusion equation, describing the scattering properties of the contact by a spatially dependent distribution matrix. The current density on the normal side of the contact ($x < 0$) is $i(x<0) = \sigma^N \partial_x f$ and consequently the total current is

$$I(x<0) = G_D^N\left(L^N \partial_x\right) f\,,$$

where f is the spatially dependent distribution matrix on the normal side in the contact. In the normal metal part the boundary condition is $f(x = -L^N) = f^N$. In a ferromagnet spin-up and spin-down states are incoherent, and hence spins non-collinear to the magnetization direction relax and only spins collinear with the magnetization will propagate sufficiently far away from the F|N-interface. It is assumed that the ferromagnet is sufficiently strong and that the contact is longer than the ferromagnetic decoherence length $\xi = \sqrt{D / h_{ex}}$, where D is the diffusion constant and hex is the exchange splitting. The decoherence length is typically very

short in ferromagnets, $\xi = 2$ nm in Ni wires.

The distribution function on the ferromagnetic side can then be represented by a two-component distribution function

$$f(x>0) = u^{\uparrow} f^{\uparrow} + u^{\downarrow} f^{\downarrow} \tag{43}$$

where $u^{\uparrow}$ and $u^{\downarrow}$ are the spin-projection matrices (34). Here a spin-accumulation collinear to the magnetization direction in the ferromagnet is taken into account. The boundary condition determined by the distribution function in the ferromagnetic part is thus

$$f^{\uparrow}(x = L^F) = f^{F\uparrow}, \quad f^{\downarrow}(x = L^F) = f^{F\downarrow} \tag{44}$$

In assumption that the resistance of the diffusive region of the contacts is much larger than the contact resistance between the normal and the ferromagnetic metal, the total current in the ferromagnet is described as

$$I(x>0) = G_D^{F\uparrow} u^{\uparrow} \partial_x f^{\uparrow} + G_D^{F\downarrow} u^{\downarrow} \partial_x f^{\downarrow} \tag{45}$$

Here, distribution function is continuous across the F|N interface, $f(0^+) = f(0^-)$. Current conservation on the left ($x<0$) and on the right ($x>0$) of the normal metal-ferromagnet interface dictates the equation, $\partial_x I = 0$, which together with the boundary conditions (44), $f(x = -L^N) = f^N$ and $f(0^+) = f(0^-)$ uniquely determine the distribution functions and hence the conductance in the diffusive contact. The current on the normal side of the contact becomes

$$eI = G_D^{\uparrow} u^{\uparrow} \left(f^F - f^N\right) u^{\downarrow} + G_D^{\downarrow} u^{\downarrow} \left(f^F - f^N\right) u^{\downarrow}$$

$$+ G_D^N \left[u^{\uparrow} \left(f^F - f^N\right) u^{\downarrow} + u^{\downarrow} \left(f^F - f^N\right) u^{\downarrow} \right] \tag{46}$$

The current in a diffusive contact thus takes the generic form (36) with $G^{\uparrow} = G_D^{\uparrow}$, $G^{\downarrow} = G_D^{\downarrow}$ and $G^{\uparrow\downarrow} = G_D^N$. The mixing conductance is thus real and only depends on the normal conductance. The latter results can be understood as a consequence of the effective spin-relaxation of spins non-collinear to the local magnetization direction. Those spins cannot propagate in the ferromagnet, and consequently the effective conductance can only depend on the conductance in the normal metal as (41) explicitly demonstrates.

In the case of the ballistic contact, the reflection and transmission coefficients appearing in (37) and (38) are diagonal in the space of the transverse channels since the transverse momentum is conserved. In a simplified model [20] the transmission channels are either closed $t = 0$ or open $t = 1$. The conductances (36), (37) can then be found by simply counting the number of propagating modes. Then the spin-dependent conductance $G_B^{\uparrow(\downarrow)} = (e^2/h) N^{\uparrow(\downarrow)}$, where $N^{\uparrow}$ ($N^{\downarrow}$) is the number of spin-up (spin-down) propagating channels The mixing conductance is determined by $G_B^{\uparrow\downarrow} = \max(G_B^{\uparrow}, G_B^{\downarrow})$ and is real. In a quantum mechanical calculation the channels just above the potential step are only partially transmitting and the channels below a potential step can have a finite transmission probability due to tunneling. Furthermore, the band structure of ferromagnetic metals is usually complicated and interband scattering exists even at ideal interfaces. The phase of the scattered wave will be relevant giving a non-vanishing imaginary part of the mixing conductance.

In the case of a tunneling contact, the transmission coefficients are exponentially small and the reflection coefficients have a magnitude close to one. The spin-dependent conductance is

$$G_T^s = \frac{e^2}{h} \sum_{nm} |t_s^{nm}|^2 \tag{47}$$

For simple models of tunnel barriers $r_s^{nm} = \delta^{nm} \exp i\phi^n = \delta r_s^{nm}$, where the phase-shift ϕ^n is spin-independent. The expansion (38) in the small correction δr_s^{nm} leads to the expression

$$\mathrm{Re}\, G_T^{\uparrow\downarrow} = \left(G_T^{\uparrow} + G_T^{\downarrow}\right)/2$$

where $G_T^{\uparrow}$ and $G_T^{\downarrow}$ are the spin-dependent tunneling conductances (47). Since the transmission coefficients in a tunnel contact are all exponentially small, the imaginary part of $G_T^{\uparrow\downarrow}$ is of the same order of magnitude as $G_T^{\uparrow}$ and $G_T^{\downarrow}$ but it is not universal and depends on the details of the contact.

3. Spin Transport in the Case of Dynamic Magnetic Field

3.1. Features of Coupling Spin Currents with Magnetic Dynamics

The interconnection between spin currents and the magnetic dynamics in F/N based magnetic multilayer nanostructures underlies the current-controlled magnetic dynamics and utilizing of the latter as new functionality in spintronic devices [21]. One is related to the *s-d* exchange interaction with localized spins and the spin-dependent scattering of spin-polarized electrons near the N|F interface. The impact of the spin current on localized spins occurs via a finite torque on the magnetic order parameter, and, vice versa, a moving magnetic order vector loses torque by emitting a spin current. The magnetic precession acts as a spin pump which transfers angular momentum from the magnetic into normal metal.

The technological potential of the mentioned magnetic nanostructures is related to utilizing transition metals (for instance, Co, Ni, Fe) that operate at ambient temperatures. Examples are current-induced tunable microwave generators (spin-torque oscillators) [22, 23], and non-volatile magnetic electronic architectures that can be randomly read, written or programmed by current pulses in a scalable manner [24]. The interaction between currents and magnetization can also cause undesirable effects such as enhanced magnetic noise in read

heads made from magnetic multilayers [25].

In the framework of the Landau-Lifshitz-Gilbert model, the impact of the spin current on the magnetic dynamics, caused by the spin transfer, reduces to change of fundamental parameters such as the gyromagnetic ratio and Gilbert damping parameters. This spin transfer is governed by the reflection and transmission matrices of the system, analogous to the scattering theory of transport and interlayer exchange coupling. In the case when the normal metal layers adjacent to the ferromagnetic layers are perfect spin sinks, the spin accumulation in the normal metal vanishes [24]. In the opposite case, the spin accumulation accompanies by the spin diffusion, which gives essential contribution to the total spin current and its interconnection with magnetic dynamics.

Spin pumping by a precessing ferromagnet is, in some sense, the reverse process of current-induced magnetization dynamics. When the pumped spin angular momentum is not quickly dissipated to the normal-metal atomic lattice, a spin accumulation builds up and creates reaction torques due to transverse-spin backflow into ferromagnets. The interplay between magnetization dynamics and the nonequilibrium spin-polarized transport in heterostructures determining magnetic properties will be considered for the case of F/N based nanostructures below.

3.2. Precession-Induced Spin Pumping Through F/N Interfaces

Characteristic properties of the precession-induced spin pumping are manifested in the model N/F/N magnetic junction schematic of which is displayed in Figure 2. The ferromagnetic layer F is a spin-dependent scatterer that governs electron transport between (left (L) and right (R)) normal metal reservoirs. The 2×2 operator $\boldsymbol{I}_l$ for the charge and spin current in l th lead ($l = L, R$) can be expressed in terms of operators $a_{\alpha m,l}(E)$ and $b_{\alpha m,l}(E)$ that annihilate a spin-α electron with energy E leaving (entering) the l th lead through the m th channel:

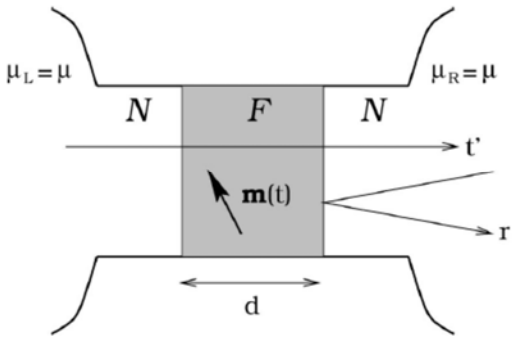

Figure 2. *Ferromagnetic layer (F) sandwiched between two normal metal layers (N). The reflection and transmission amplitudes r and t' govern the spin current pumped into the right lead.*

$$I_l^{\alpha\beta}(t) = \frac{e}{h}\sum_m \int dE dE' e^{i(E-E')t/\hbar} \times \left[a^{\dagger}_{\beta m,l}(E) a_{\alpha m,l}(E') - b^{\dagger}_{\beta m,l}(E) b_{\alpha m,l}(E') \right] \quad (48)$$

When the scattering matrix $s^{\alpha\beta}_{mn,ll'}(t)$ of the ferromagnetic layer varies slowly on the time scales of electronic relaxation in the system, an adiabatic approximation may be used. The annihilation operators for particles entering the reservoirs are then related to the operators of the outgoing states by the instantaneous value of the scattering matrix $b_{\alpha m,l}(E) = s^{\alpha\beta}_{mn,ll'}(t) a_{\beta n,l'}(E)$. In terms of $a_{\alpha m,l}$ only, we can evaluate the expectation value $\left\langle I_l^{\alpha\beta}(t) \right\rangle$ of the current operator using $\left\langle a^{\dagger}_{\alpha m,l}(E) a_{\beta n,l'}(E') \right\rangle = f_l(E)\delta_{\alpha,\beta}\delta_{m,n}\delta_{l,l'}\delta(E-E')$, where $f_l(E)$ is the (isotropic distribution function in the l th reservoir. When the scattering matrix depends on a single time-dependent parameter $X(t)$, then the Fourier transform of the current expectation value $I_l(\omega) = \int dt e^{i\omega t} I_l(t)$ can be written as

$$I_l(\omega) = g_{X,l}(\omega) X(\omega) \quad (49)$$

In terms of a frequency ω- and X-dependent parameter $g_{X,l}$:

$$g_{X,l}(\omega) = -\frac{e\omega}{4\pi}\sum_{l'}\int dE \left(-\frac{\partial f_{l'}(E)}{\partial E} \right) \times \sum_{mn} \left(-\frac{\partial s_{mn,ll'}(E)}{\partial X} s^{\dagger}_{mn,ll'}(E) - \text{H.c.} \right) \quad (50)$$

Equation (49) is the first-order (in frequency) correction to the dc Landauer-Büttiker formula [11]. At equilibrium $f_R(E) = f_L(E)$, (49) is the lowest-order nonvanishing contribution to the current. Furthermore, at sufficiently low temperatures, we can approximate $\partial f_l(E)/\partial E$ by a δ-function centered at Fermi energy. The expectation value of the 2×2 particle-number operator $Q_l(\omega)$ (defined by $I_l(t) = dQ_l(t)/dt$) in time or by $I_l(\omega) = -i\omega Q_l(\omega)$ in frequency domain) for the l th reservoir is then given by

$$Q_l(\omega) = \left(\frac{e}{4\pi i} \sum_{mnl'} \frac{\partial s_{mn,ll'}(E)}{\partial X} s^{\dagger}_{mn,ll'}(E) + \text{H.c.} \right) X(\omega) \quad (51)$$

where the scattering matrices are evaluated at the Fermi energy. Because the prefactor on the right-hand side of (51) does not depend on frequency ω, the equation is also valid in time domain. The change in particle number $\delta Q_l(t)$ is proportional to the modulation $\delta X(t)$ of parameter X and the 2×2 matrix current (directed into the normal-metal leads) reads

$$I_l(t) = e\frac{dn_l}{dX}\frac{dX(t)}{dt} \quad (52)$$

where the "matrix emissivity" into lead l is

$$\frac{\partial n_l}{\partial X} = \frac{1}{4\pi i}\sum_{mnl'} \frac{\partial s_{mn,ll'}(E)}{\partial X} s^{\dagger}_{mn,ll'}(E) + \text{H.c.} \tag{53}$$

If the spin-flip scattering in the ferromagnetic layer is disregarded, the scattering matrix s can be written in terms of the spin-up and spin-down scattering coefficients $s^{\uparrow(\downarrow)}$ using the projection matrices $u^{\uparrow} = (1+\boldsymbol{\sigma}\cdot\boldsymbol{m})/2$ and $u^{\downarrow} = (1-\boldsymbol{\sigma}\cdot\boldsymbol{m})/2$:

$$s_{mn,ll'} = s^{\uparrow}_{mn,ll'}u^{\uparrow} + s^{\downarrow}_{mn,ll'}u^{\downarrow} \tag{54}$$

The spin current pumped by the magnetization precession is obtained by identifying $X(t) = \varphi(t)$, where φ is the azimuthal angle of the magnetization direction in the plane perpendicular to the precession axis. For simplicity, we assume that the magnetization rotates around the yax is: $\boldsymbol{m} = (\sin\varphi, 0, \cos\varphi)$. Using (54), it is then easy to calculate the emissivity (53) for this process:

$$\frac{\partial n_l}{\partial \varphi} = -\frac{1}{4\pi}\left[A_r\sigma_y + A_i(\sigma_x\cos\varphi - \sigma_y\sin\varphi)\right] \tag{55}$$

where $A_r(A_i) = \text{Re(Im)}[\text{g}^{\uparrow\downarrow} - t^{\uparrow\downarrow}]$. Expanding the 2×2 current into isotropic and traceless components,

$$I = \frac{1}{2}I_c - \frac{e}{\hbar}\boldsymbol{\sigma}\cdot\boldsymbol{I}_s \tag{56}$$

here the charge current I_c and spin current $\boldsymbol{I}_s$ are identified. Due to equations (52), (55) and (56) the charge current vanishes, $I_c = 0$, and the spin current

$$\boldsymbol{I}_s = \left(A_i\cos\varphi, A_r - A_i\sin\varphi\right)\frac{\hbar}{4\pi}\frac{d\varphi}{dt} \tag{57}$$

can be rewritten as

$$\boldsymbol{I}_s^{\text{pump}} = \frac{\hbar}{4\pi}\left(A_r\boldsymbol{m}\times\frac{d\boldsymbol{m}}{dt} - A_i\frac{d\boldsymbol{m}}{dt}\right) \tag{58}$$

This current into a given N layer depends on the complex-valued parameter $A = A_r + iA_i$ (the "spin-pumping conductance") and the time-dependent order parameter of the ferromagnet $\boldsymbol{m}(t)$. In addition, $A = g^{\uparrow\downarrow} - t^{\uparrow\downarrow}$ depends on the scattering matrix of the ferromagnetic layer since

$$g^{\sigma\sigma'} = \sum_{mn}\left[\delta_{mn} - r^{\sigma}_{mn}(\mathrm{r}^{\sigma'}_{mn})^{*}\right] \tag{59}$$

is the dimensionless dc conductance matrix [10, 26] and

$$t^{\uparrow\downarrow} = \sum_{mn} t'^{\uparrow}_{mn}(t'^{\downarrow}_{mn})^{*} \tag{60}$$

Here $r^{\uparrow}_{mn}$ ($r^{\downarrow}_{mn}$) is a reflection coefficient for spin-up ~spin-down! electrons on the normal-metal side and $t'^{\uparrow(\downarrow)}_{mn}$ is a transmission coefficient for spin-up (spin-down) electrons across the ferromagnetic film from the opposite reservoir into the normal-metal layer, where *m* and *n* label the transverse modes at the Fermi energy in the normal metal films. The magnetization can take arbitrary directions, in particular, $\boldsymbol{m}(t)$ may be far away from its equilibrium value. In such a case, the scattering matrix itself can depend on the orientation of the magnetization, and one has to use in (58).

When the ferromagnetic film is thicker than its transverse spin-coherence length $d > \pi/(k^{\uparrow} - k^{\downarrow})$, where $k_F^{\uparrow\downarrow}$ are the spin-dependent Fermi wave vectors, $t^{\uparrow\downarrow}$ vanishes [26], the spin pumping through a given *F-N* interface is governed entirely by the interfacial mixing conductance $A = g^{\uparrow\downarrow} = g_r^{\uparrow\downarrow} + ig_i^{\uparrow\downarrow}$, and we can consider only one of the two interfaces.

The spin current (58) leads to a damping of the ferromagnetic precession, resulting in a faster alignment of the magnetization with the (effective) applied magnetic field $\boldsymbol{H}_{eff}$. The pumped spins are entirely absorbed by the attached ideal reservoirs. Thereto the enhancement rate of damping is accompanied by an energy flow out of the ferromagnet, until a steady-state is established in the companied *F/N* system. For simplicity, assume a magnetization which at time t starts rotating around the vector of the magnetic field, $\boldsymbol{m}(t) \perp \boldsymbol{H}_{eff}$. In short interval of time δt, it slowly changes toa short interval of time δt, it slowly (i.e., adiabatically) changes to $\boldsymbol{m}(t+\delta t) = \boldsymbol{m}(t) + \delta\boldsymbol{m}$. In the presence of a large but finitenonmagnetic reservoir without any spin-flip scattering attached to one side of the ferromagnet, this process can be expected to induce a nonvanishing spin accumulation

$$\mu_s = \int d\varepsilon \,\text{Tr}[\sigma \boldsymbol{f}(\varepsilon)] \tag{61}$$

where σ is the Pauli matrix vector and $\boldsymbol{f}(\varepsilon)$ is the 2×2 matrix distribution function at a given energy ε. For a slow enough variation of $\boldsymbol{m}(t)$, this nonequilibrium spin imbalance must flow back into the ferromagnet, canceling any spin current generated by the magnetization rotation, since, due to the adiabatic assumption, the system is always in a steady state.

For the spins accumulated in the reservoir along the magnetic field, $\boldsymbol{\mu}_s \parallel \boldsymbol{H}_{eff}$ flow of N_s spins into the normal metal transfers the energy $\Delta E_N = N_s\mu_s/2$ and angular-momentum $\Delta L_N = N_s\mu_s/2$ (directed along $\boldsymbol{H}_{eff}$). By the conservation laws, $\Delta E_N = N_s\mu_s/2$ and $\Delta L_F = -\Delta L_N$, for the corresponding values in the ferromagnet. Using the magnetic energy $\Delta E_f = -\gamma\Delta L_F H_{eff}$, where γ is the gyromagnetic ratio of the ferromagnet, we than find that $\mu_s = \hbar\gamma H_{eff} = \hbar\omega$, where ω is the Larmor frequency of precession in the effective field. The spin-up and spin-down chemical potentials in the normal metal split by μ_s, the energy corresponding to the frequency of the perturbation.

The above mentioned the backflow of spin current $\boldsymbol{I}_s^{\text{back}}$, which equal to the pumping current $\boldsymbol{I}_s^{\text{pump}}$ described by the

expression

$$\boldsymbol{I}_s^{\text{back}} = \frac{1}{2\pi}\left(g_r^{\uparrow\downarrow}\boldsymbol{\mu}_s + g_i^{\uparrow\downarrow}\boldsymbol{m}\times\boldsymbol{\mu}_s\right)$$

$$= \frac{\hbar}{4\pi}\left(g_r^{\uparrow\downarrow}\boldsymbol{m}\times\frac{d\boldsymbol{m}}{dt} - g_i^{\uparrow\downarrow}\frac{d\boldsymbol{m}}{dt}\right) \qquad (62)$$

Here, it is used that $\mu_s = \hbar\omega$ and $\boldsymbol{\mu}_s \perp \boldsymbol{m}$, since by the conservation of angular momentum, the spin transfer is proportional to the change in the direction $\delta\boldsymbol{m} \perp \boldsymbol{m}$. Thus for the case of a single and finite reservoir (58) is recovered. It is easy to repeat the proof for an arbitrary initial alignment of $\boldsymbol{m}(t)$ with $\boldsymbol{H}_{eff}$. The expressions for the adiabatic spin pumping are not the whole story, since spin-flip scattering is essential. In this case, the spin build-up occurs in the normal metal at dynamic equilibrium. Then, the contribution to $\boldsymbol{I}_s$ due to the spin-accumulation-driven current $\boldsymbol{I}_s^{\text{back}}$ back into the ferromagnet:

$$\boldsymbol{I}_s = \boldsymbol{I}_s^{\text{pump}} - \boldsymbol{I}_s^{\text{back}} \qquad (63)$$

which vanishes in the absence of spin-flip scattering.

The spin current out of the ferromagnet carries angular momentum perpendicular to the magnetization direction. By conservation of angular momentum, the spins ejected by $\boldsymbol{I}_s$ correspond to a torque $\boldsymbol{T} = -\boldsymbol{I}_s$ on the ferromagnet. If possible, interfacial spin-flip processes are disregarded, the torque t is entirely transferred to the coherent magnetization precession. The dynamics of the ferromagnet can then be described by a generalized Landau-Lifshitz-Gilbert (LLG) equation [5]

$$\frac{d\boldsymbol{m}}{dt} = -\gamma\boldsymbol{m}\times\boldsymbol{H}_{eff} + \alpha_0\boldsymbol{m}\frac{d\boldsymbol{m}}{dt} + \frac{\gamma}{M_s V}\boldsymbol{I}_s \qquad (64)$$

where α_0 is the dimensionless bulk Gilbert damping constant, M_s is the saturation magnetization of the ferromagnet, and V is its volume. The intrinsic bulk constant α_0 is smaller than the total Gilbert damping $\alpha = \alpha_0 + \alpha'$. The additional damping α' caused by the spin pumping is observable in, for example, FMR spectra here.

3.3. Spin-Accumulation-Driven Backflow in the F/N Bilayer

The precession of the magnetization does not cause any charge current in the system. The spin accumulation or nonequilibrium chemical potential imbalance $\boldsymbol{\mu}_s$ (similar to (61), but spatially dependent now) in the normal metal is a vector, which depends on the distance from the interface *x*, $0 < x < L$, where *L* is the thickness of the normal-metal film (Figure 3). When the ferromagnetic magnetization steadily rotates around the *z* axis, $\boldsymbol{m}\times\partial_t\boldsymbol{m}$ and the normal-metal spin accumulation $\boldsymbol{\mu}_s(x)$ are oriented along *z*, as depicted in Figure 2. There is no spin imbalance in the ferromagnet, because $\boldsymbol{\mu}_s(x)$ is perpendicular to the magnetization direction $\boldsymbol{m}$. As shown below, the time-dependent $\boldsymbol{m}_s$ is also perpendicular to $\boldsymbol{m}$ even in the case of a precessing ferromagnet with time-dependent instantaneous rotation axis, as long as the precession frequency ω is smaller than the spin-flip rate τ_{sf}^{-1} in the normal metal.

The spin accumulation diffuses into the normal metal as

$$i\omega\boldsymbol{\mu}_s = D\partial_x^2\boldsymbol{\mu}_s - \tau_{sf}^{-1}\boldsymbol{\mu}_s \qquad (65)$$

where *D* is the diffusion coefficient. The boundary conditions are determined by the continuity of the spin current from the ferromagnet into the normal metal at $x = 0$ and the vanishing of the spin current at the outer boundary $x = L$:

$$x = 0:\ \partial_x\boldsymbol{\mu}_S = -2(\hbar NSD)^{-1}\boldsymbol{I}_s,$$

$$x = L:\ \partial_x\boldsymbol{\mu}_s = 0 \qquad (66)$$

where N is the (one-spin) density of states in the lalyer and *S* is the area of the interface. The solution to (65) with the boundary conditions (66) is

$$\boldsymbol{\mu}_s(x) = \frac{\cosh k(x-L)}{\sinh kL}\frac{2\boldsymbol{I}_s}{\hbar NSDk} \qquad (67)$$

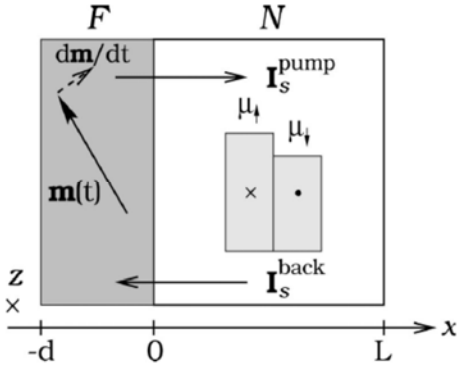

Figure 3. *Schematic of the F/N bilayer. Precession of the magnetization direction* $\boldsymbol{m}(t)$ *of the ferromagnet F pumps spin into the adjacent normal-metal layer N by inducing a spin current* $\boldsymbol{I}_s^{\text{pump}}$. *This lead to a build-up of the normal metal spin accumulation which either relaxes by spin-flip scattering or flows back into the ferromagnet as* $\boldsymbol{I}_s^{\text{back}}$.

with the wave vector $k = \lambda_{sd}^{-1}\sqrt{1 + i\omega\tau_{sf}}$, where $\lambda_{sd} = \sqrt{D\tau_{sf}}$ is the spin-flip diffusion length in the normal metal.

Using relation $D = v_F^2\tau_{el}/3$ between the diffusion coefficient *D*, the Fermi velocity v_F, and the elastic scattering time τ_{el}, we find for the spin-diffusion length $\lambda_{sd} = v_F\sqrt{\tau_{el}\tau_{sf}/3}$. An effective energy-level spacing of the states participating in the spin-flip scattering events in a thick layer can be defined by $\delta_{sd} = (NS\lambda_{sd})^{-1}$. The spin-accumulation-driven spin current $\boldsymbol{I}_s^{\text{back}}$ through the interface reads

$$\boldsymbol{I}_s^{\text{back}} = \frac{1}{8\pi}\Big[2g_r^{\uparrow\downarrow}\boldsymbol{\mu}_s(x=0) + 2g_i^{\uparrow\downarrow}\boldsymbol{m}\times\boldsymbol{\mu}_s(x=0) + (g^{\uparrow\uparrow} + g^{\downarrow\downarrow} - 2g_r^{\uparrow\downarrow})\big(\boldsymbol{m}\cdot\boldsymbol{\mu}_s(x=0)\big)\Big] \quad (68)$$

Substitution of (67) into (68) gives total spin current

$$\boldsymbol{I}_s = \boldsymbol{I}_s^{\text{pump}} - \frac{\beta}{2}\Big[2g_r^{\uparrow\downarrow}\boldsymbol{I}_s + 2g_i^{\uparrow\downarrow}\boldsymbol{m}\times\boldsymbol{I}_s + (g^{\uparrow\uparrow} + g^{\downarrow\downarrow} - 2g_r^{\uparrow\downarrow})\big(\boldsymbol{m}\cdot\boldsymbol{I}_s\big)\boldsymbol{m}\Big] \quad (69)$$

where the spin current returning into the ferromagnet is governed by the backflow factor $\beta = (\tau_{sf}\delta_{\text{sd}})/(\hbar\tanh(L/\lambda_{sd}))$. When the normal metals is shorter than the spin-diffusion length ($L \ll \lambda_{sd}$), $\beta \to \tau_{sf}\delta/h$, where $\delta = (NSL)^{-1}$. Basically, β is therefore the ratio between the energy level spacing of the normal-metal layer with thickness $L_{sf} = \min(L, \lambda_{sd})$ and the spin-flip rate.

By inverting (69) the total spin current $\boldsymbol{I}_s$ can be express in terms of the pumped spin current $\boldsymbol{I}_s^{\text{pump}}$,

$$\boldsymbol{I}_s = \left[1 + 2g_r^{\uparrow\downarrow} + \frac{(\beta g_i^{\uparrow\downarrow})^2}{1+\beta g_r^{\uparrow\downarrow}}\right]^{-1}\left(1 - \frac{\beta g_i^{\uparrow\downarrow}}{1+\beta g_r^{\uparrow\downarrow}}\boldsymbol{m}\times\right)\boldsymbol{I}_s^{\text{pump}} \quad (70)$$

Substitution (58) into (70) results in total spin current $\boldsymbol{I}_s$, which is described by the equation of the form (58) but with a redefined spin-pumping conductance $A' = A'_r + A'_i$, i.e.

$$\boldsymbol{I}_s = \frac{\hbar}{4\pi}\left(A'_r\boldsymbol{m}\times\frac{d\boldsymbol{m}}{dt} - A'_i\frac{d\boldsymbol{m}}{dt}\right) \quad (71)$$

Here A' is function of the mixing conductance $g^{\uparrow\downarrow}$ and the backflow factor β, $A' = A'(g^{\uparrow\downarrow}, \beta)$. For realistic $F|N$ interfaces $g_i^{\uparrow\downarrow} \ll g_r^{\uparrow\downarrow}$, so that $g^{\uparrow\downarrow} \approx g_r^{\uparrow\downarrow}$ and, consequently, A'_i vanishes. After substituting (71) into (74) renormalizes its Gilbert dumping constant, α_0, so that $\alpha_0 \to \alpha = \alpha_0 + \alpha'$, where

$$\alpha' = \frac{g_L g^{\uparrow\downarrow}}{4\pi\mu}\left[1 + g^{\uparrow\downarrow}\frac{(\tau_{sf}\delta_{sd}/h\beta g_i^{\uparrow\downarrow})^2}{\tanh(L/\lambda_{sd})}\right]^{-1} \quad (72)$$

is the additional damping constant due to the interfacial F|Ncoupling. Here, g_L is the g factor and μ is the total layer magnetic moment in units of μ_{B}. When $L \to \infty$, (72) reduces to a simple result: $\alpha' = g_L g_{eff}^{\uparrow\downarrow}/4\pi$ where $1/g_{eff}^{\uparrow\downarrow} = 1/g^{\uparrow\downarrow} + R_{sd}$, where $R_{sd} = \tau_{sf}\times\delta_{sd}/h$ is the resistance (per spin, in units of h/e^2) of the normal-metal layer of thickness λ_{sd}. It follows that the effective spin pumping out of the ferromagnet is governed by $g_{eff}^{\uparrow\downarrow}$, i.e., the conductance of the $F|N$ interface in series with diffusive normal-metal film with thickness λ_{sd}.

The second factor on the right-hand side of (72) suppresses the additional Gilbert damping due to the spin angular momentum that diffuses back into the ferromagnet. Because spins accumulate in the normal metal perpendicular to the ferromagnetic magnetization, the spin–accumulation-driven transport across the F|Ncontact, as well as the spin pumping, is governed by a mixing conductance.

4. Conclusions

The spin transport In the F/N based magnetic nanostructures in magneto-static and magneto-dynamic cases have been studied in the framework of the modified Stoner model. Using the modified quantum-kinetic equation for the non-linear Green functions and the spin-dependent scattering matrix, the spin currents through and near the F|N interface are described. In the magneto-static case, the parametric dependence of the spin current on the relative orientation of the spin polarization and magnetization is shown. In the magneto-dynamic case of the magnetization precession, the precession-induced spin pumping into the normal metal layer is described. The accompanying effect of the spin accumulation and the spin backflow exerting via the spin torque on the magnetization precession are considered.

References

[1] Equation Chapter 1 Section 1 L. Zutic, J. Fabian, and S. Sarma, "Spintronics: Fundamentals and applications," Rev. Mod. Phys., vol. 76, pp. 323-413, April 2004.

[2] J. Fisher, O. Gomonay, R. Schlitz, K. Ganzhorn, N. Vliestra, M. Althammer et al. "Spin Hall magnetoresistance in antiferromagnet/normal metal heterostructures," arXiv:cond-mat.mes-hall, September 2017.

[3] Y. Tserkovnyak, A. Brataas, G. E. Bauer, and B. I. Halperin, "Nonlocal magnetization dynamics in ferromagnetic heterostructures," Rev. Mod. Phys., vol. 77, pp. 1375-1421, October 2005.

[4] A. Brataas, Yu. V. Nazarov, and G. E. W. Bauer, "Spin-transport in multi-terminal normal metal – ferromagnet systems with non-collinear magnetization." Eur. Phys. J., vol. B22, pp. 99-110, February 2001.

[5] J. C. Slonczewski, "Current-driven excitation of magnetic multilayers," J. Magn. Magn. Mater.. vol. 159, pp L1-L7, February 1996.

[6] A. Manchona, N. Strelkova, B. N. Ryzhanovaa, B. A. Vedyayeva, B. B. Dienya, J. C. Slonczewski, "Theoretical investigation of the relationship between spin torque and magnetoresistance in spin-valves and magnetic tunnel junctions," J. Magn. Magn. Mater.. vol. 316, pp L977-L979, March 2007.

[7] J. Akerman, "Toward a Universal Memory," Science, vol. 308, pp. 508–510, April 2005.

[8] J. Katine, and E. F. Fullerton, "Device implications of spin-transfer torques," J. Magn. Magn. Mater., 320, 1217–1226, April 2008.

[9] G. Binasch, P. Grünberg, Saurenbach F., and W. Ziman, "Enhanced magnetoresistance in layered magnetic structures with antiferromagnetic interlayer exchange," Phys. Rev., vol. B 39, pp. 4828-4830, March 1989.

[10] X. Waintal, E. B. Myers, P. W. Brouwer, and D. C. Ralph, "Role of spin-dependent interface scattering in generating current-induced torques in magnetic multilayers," Phys. Rev., vol. B 62, pp. 12317-12327, November 2000.

[11] M. Büttiker, "Four-Terminal Phase-Coherent Conductance," Phys. Rev. Lett., vol. 57, pp. 1761-1764, October 1986.

[12] P. Danielewicz, "Quantum Theory of Nonequilibrium Processes," Ann. Phys., vol. 152**,** pp. 234-304, November 1984.

[13] P. Myöhänen, A. Stan, G. Stefanucci and R. van Leeuwen. "Kadanoff-Baym approach to quantum transport through interacting nanoscale systems: From the transient to the steady-state regime," Phys. Rev. vol. B 80, pp. 115107-1─115107-15, September 2009.

[14] R. Cheng, J.-J. Zhu, and D. Xiao, "Dynamic Feedback in Ferromagnet/Spin Hall Heterostructures," Phys. Rev. B vol. 117, pp. 097202-097206, August 2016.

[15] A. Manchon, H. C. Koo, J. Nitta, S. M. Frolov, and R. A. Duine, " New Perspective for Rashba Spin-Orbit Coupling," Nature Materials, vol. 36, pp. 871-382, August 2015.

[16] I. M. Miron, G Gilles, S. Auffret, B. Rodmacq, A. Schuhl, S. Pizzini, J. Vogel, and P. Gambardella, "Current-driven spin torque induced by the Rashba effect in a ferromagnetic metal layer," Nat. Matter., vol. 9, pp. 230-234, January 2010.

[17] A. Brataas, Y. V. Nazarov, J. Inoue, G. E. W. Bauer, "Spin accumulation in small ferromagnetic doublebarrier junctions," Phys. Rev. vol., B 59, 93-96, January 1999.

[18] D. Huertas Hernando, A. Brataas, Y. V. Nazarov, G. E. W. Bauer, "Conductance modulation by spin precessing in noncollinear ferromagnet/normal metal ferromagnetic multilayers," Phys. Rev., vol. B 62, pp. 5700-5712, September 2000.

[19] M. A. M. Gijos, and G. E. W. Bauer, "Perpendicular giant magnetoresistance of magnetic multilayers," Adv. Phys., vol. 46, pp. 285-445, February 1997.

[20] G. E. W. Bauer, "Perpendicular transport through magnetic multilayers," Phys. Rev. Lett., vol. 69, pp. 1676-1679, September 1992.

[21] S. D. Bader and S. S. P. Parkin, "Spintronics."Annual Review of Condensed Matter Physics, vol. 1, pp. 71-88, April 2010.

[22] T. J. Silva and W. H. Rippard, "Developments in nano-oscillators based upon spin-transfer point-contact devices," J. Magn. Magn. Mater. vol. 320, pp. 1260-1271, April 2008.

[23] P M Braganca, B A Gurney, B A Wilson, J A Katine, S Maat, and J R Childress, "Nanoscale magnetic field detection using a spin torque oscillator," Nanotechnology, vol. 21, pp. 235202-1─235202-6, May 2010.

[24] S. Matsunaga, K. Hiyama, A. Matsumoto, S. Ikeda, H. Hasegawa, K. Miura, J. Hayakawa, T. Endoh, H. Ohno, and T. Hanyu, "Standby-power-free compact ternary content-addressable memory cell chip using magnetic tunnel junction devices," Applied Physics Express, vol. 2, pp. 023004-1─023004-6, February 2009.

[25] K. Nagasaka, "CPP-GMR technology for magnetic read heads of future high-density recording systems," J. Magn. Magn. Mater., vol. 321, pp. 508-51, June 2009.

[26] M. D. Stiles and A. Zangwill, "Anatomy of spin-transfer torque," Phys. Rev., vol. B 66, pp. 014407-1─014407-14, June 2002.