ISSN 1064-2269, Journal of Communications Technology and Electronics, 2019, Vol. 64, No. 12, pp. 1422–1430. © Pleiades Publishing, Inc., 2019.
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NANOELECTRONICS

Spin Polarization of Nonequilibrium Conduction Electrons in Magnetic Junctions

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Received March 14, 2019; revised May 14, 2019; accepted May 15, 2019

Abstract—An original equation for nonequilibrium spin polarization in magnetic junctions based on dynamic equations for magnetic moment was obtained. Spatial nonuniformity of the distribution of carriers is taken into account in the equations, and the magnetic moment is averaged over the ensemble of nonequilibrium spin-injected electrons. The solution to the dynamic equations for the magnetic moment is used to estimate the probability of spin-flip electron transitions upon spin injection, and the solution to the equation for the nonequilibrium spin polarization in magnetic junctions is used to calculate frequencies of photon emission or absorption at energies corresponding to the energy of effective exchange splitting of spin subbands.

DOI: 10.1134/S1064226919100139

INTRODUCTION

Magnetic heterostructures with spin-polarized current flow are employed in terahertz spintronics, 3D nanomagnetic devices, sensors, and storage devices [1–16]. Spin polarization of the current that flows in a heterostructure may lead to significant nonequilibrium accumulation of spins in a finite volume of a ferromagnetic material. The spin relaxation of nonequilibrium electrons may occur via interaction with ferromagnetic lattice with excitation of oscillations of magnetization or magnetization switching [17, 18]. Note less data on the interaction of spins of conduction electrons and magnetic lattice of a ferromagnetic material under resonance conditions for relaxation processes (including photoemission) caused by spin flip of conduction (*s*) electrons [19–24]. Spin-relaxation transitions of conduction electrons between spin subbands of a ferromagnetic material are indirectly stimulated via the *sd*-exchange interaction with electromagnetic wave. Such spin-flip electron transitions can be accompanied by emission or absorption of photons [25] with the frequency determined by the energy of effective exchange splitting of the spin subbands. For several transitions, the splitting energy of the spin subbands corresponds to the photon energies of the terahertz frequency range [26–32]. In this case, the heterostructures can be employed in technical applications as the basic units of terahertz radiation sources including devices working at room temperature.

Conduction electrons may interact with external electromagnetic radiation via the *sd*-exchange with vector potential $\vec{A}$ of external field included in the energy of the exchange *sd*-interaction [24]. Emission via such a relativistic correction to the momentum of conduction electrons (*s*-electrons) modulates the *sd*-interaction and may cause radiative transitions of *s*-electrons with spin flip. Such an interaction channel is more efficient than conventional multipole channels by several orders of magnitude [33]. Its high efficiency has been demonstrated in [24] using estimations of probabilities of spin-flip electronic transitions. Note that the effect of the dynamics of nonequilibrium population of electron subbands and the effect of thermostat of electron system on the probability of energy transitions of electrons with spin flip have note been taken into account. For a more correct analysis of the photon emission, a model of the spin-flip transitions based on the dynamic equations for magnetic moment has been proposed in [31, 32]. Such equations determine the time dependence of the averaged magnetic moment of nonequilibrium injected electrons and the spatial nonuniformity of the distribution of carriers. In this work, we use the dynamic equations for magnetic moment of [31, 32] to derive an equation for nonequilibrium spin polarization in magnetic junctions. The solution to such an equation makes it possible to calculate the frequencies of photon emission and absorption for indirect energy transitions of electrons with spin flip.

1. MODEL OF MAGNETIC JUNCTION

We consider a scheme of a typical magnetic junction that is similar to the scheme of [29] (Fig. 1). Here, the nonequilibrium spin injection in a ferromagnetic material is achieved using effective spin injection in the presence of the current flow.

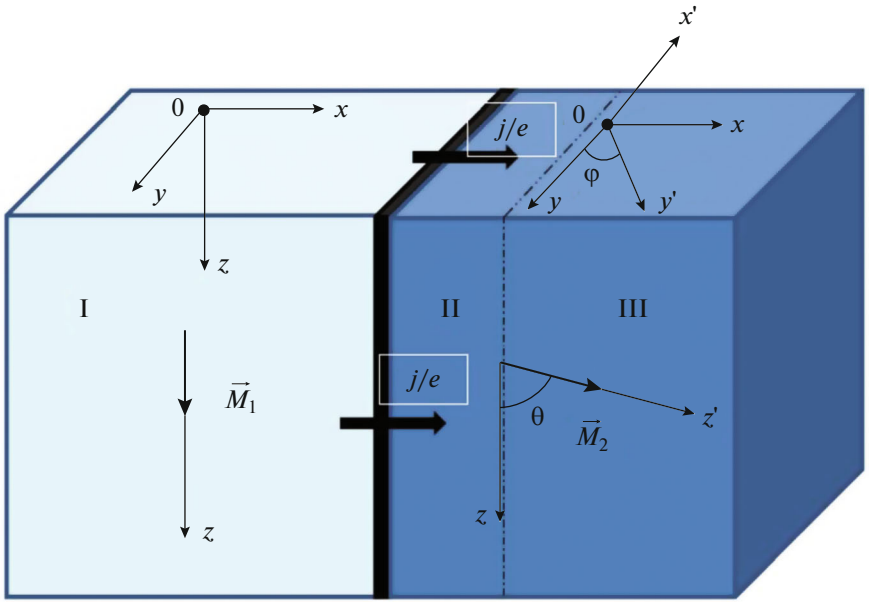

Fig. 1. Scheme of spin-injection magnetic junction: (I) first ferromagnetic material that serves as injector, (II) injection region of the second ferromagnetic material, and (III) fragment of the second ferromagnetic material outside the injection region.

In the derivation of the original equations, we take into account the dependence of the nonequilibrium spin polarization on the x, y, and z coordinates [29] and assume that the spin polarization is changed from the initial state due to the presence of an injection current pulse. A current pulse flows through the interface of metal ferromagnetic materials I and II magnetization vectors of which $\vec{M}_1$ and $\vec{M}_2$ are parallel to quantization axes in materials I and II, and θ is the angle between vectors $\vec{M}_1$ and $\vec{M}_2$ (Fig. 1). The exchange interaction of ferromagnetic materials I and II is eliminated owing to the presence of an ultrathin buffer layer made of nonmagnetic insulator or metal. In the presence of the current flow through the spin-injection junction from the first ferromagnetic material (region I), uncompensated spins are injected to region II of the second ferromagnetic material (below, working region of the magnetic junction). Then, the injected carriers pass through region III that serves as a collector.

Figure 2 shows changes of the quasi-Fermi levels for mean populations of the spin subbands in regions I and II (and the energy relaxation transitions of electrons with spin flip in region II) [29]. For the terahertz frequency range, we may assume relatively weak oscillations of the magnetization of ferromagnetic lattices (typical eigenfrequencies of such oscillations are 10^{10}–10^{11} Hz). Thus, the magnetic lattice is static in the first approximation. Note also that the oscillations and magnetic switching of the ferromagnetic lattice due to transfer of torque from the s electrons are observed in magnetic layers with thicknesses of about several nanometers. We assume that the thickness of the layer in which the spin state remains unchanged (working layer) is significantly greater than the above thickness. Thus, the state with stationary direction of magnetization of the ferromagnetic material becomes preferred with respect to energy. In this case, the oscillation frequencies of the spin subsystem of the s electrons are significantly greater than the frequencies of the $3d$ electrons that are localized on atoms of crystal lattice and related to its magnetization. The above facts make it possible to disregard the dynamics of magnetic lattice at the terahertz frequencies.

Then, we assume that the conduction electrons have isotropic parabolic energy spectrum with a certain effective mass. We also assume that such electrons are affected by the exchange interaction with the d electrons (sd exchange), external electromagnetic field, and external equilibrium system of the thermostat. The exchange interaction of the s electrons that determine the conductivity of a ferromagnetic material with the d electrons that are localized on the atoms of the crystal lattice depends on the spin state of the s and d electrons on the assumption that the spin state of the d electrons is related to the magnetization of the crystal lattice of the ferromagnetic material. Such interaction is described using the sd exchange process [34]. In the analysis of the mean populations of injected electrons in the spin subbands, we may introduce quasi-stationary Fermi levels in each subband as a correction to the equilibrium position of the Fermi level. In addition, we take into account the condition for electroneutrality in such a way that the total number of electrons in the working region of the magnetic junction remains unchanged. Thus, we obtain different positions of the quasi-levels in the spin subbands, which may lead to a negative spin temperature (that

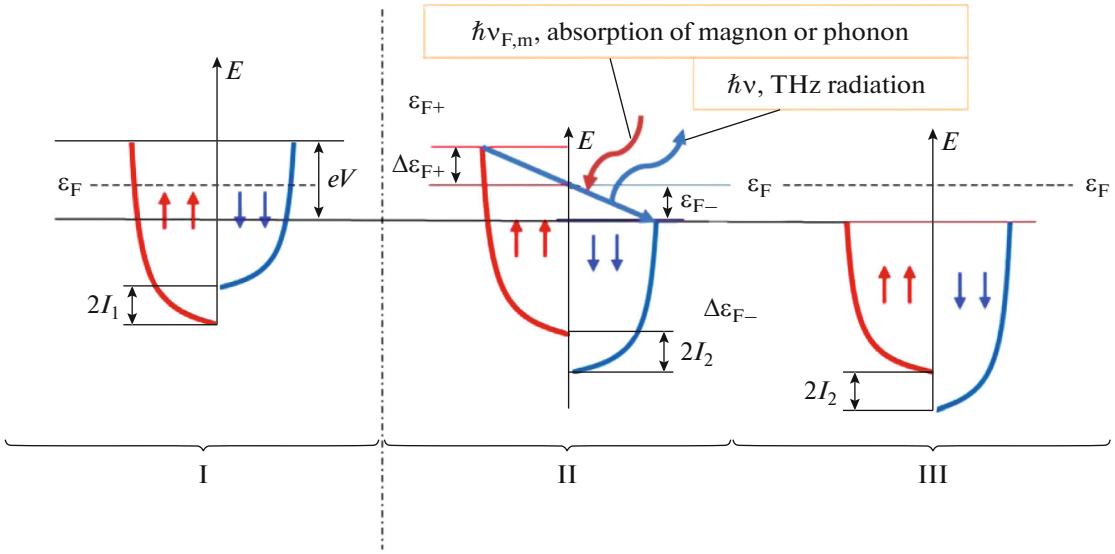

Fig. 2. Schemes of electronic energy levels and quasi-Fermi levels (I) in the first ferromagnetic material, (II) for nonequilibrium electrons in the second ferromagnetic material, and (III) for equilibrium electrons in the second ferromagnetic material outside the injection region.

has been predicted in [35]) upon current-induced spin injection in the magnetic junction.

2. HAMILTONIAN OF INJECTED ELECTRONS

The Hamiltonian of a single injected electron with allowance for the interaction with electromagnetic wave for the Pauli equation can be represented as [24, 31, 32]

$$\hat{H} = \hat{H}_\mathrm{k} + \hat{H}_{sd} + \hat{H}_A. \quad (1)$$

Here $\hat{H}_\mathrm{k}$ is the operator of kinetic energy of electron that is equal to a product of $\vec{p}^2/2m^*$ and unity matrix 2×2, $\hat{H}_{sd}$ is the operator of the *sd* exchange interaction that is generally given by

$$\hat{H}_{sd} = -\vec{\hat{\sigma}}\mu_\mathrm{B}\vec{H}^0_{sd}, \quad (2)$$

where the vector matrix is written in the following way using the unit vectors of coordinate axes: $\vec{\hat{\sigma}} = \vec{x}\hat{\sigma}_x + \vec{y}\hat{\sigma}_y + \vec{z}\hat{\sigma}_z$, $\vec{H}^0_{sd}$ is the intrinsic effective *sd*-exchange field, μ_B is the Bohr magneton, and $\hat{H}_A$ is the operator of the interaction with electromagnetic wave with vector potential $\vec{A}$ that is given by [24]

$$\hat{H}_A = -\vec{\hat{\sigma}}\mu_\mathrm{B}\vec{H}_A, \quad (3)$$

where the *sd*-exchange field induced by the electromagnetic wave is represented as

$$\vec{H}_A = \frac{e}{2c}\left(\frac{\partial\vec{H}^0_{sd}}{\partial\vec{\hat{p}}}\vec{A} + \vec{A}\frac{\partial\vec{H}^0_{sd}}{\partial\vec{\hat{p}}}\right). \quad (4)$$

We use the first two terms of Hamiltonian (1) to obtain the spinor solution with an up or down spin corresponding to two parabolic spin subbands with energy splitting by $2\mu_\mathrm{B}\left|\vec{H}^0_{sd}\right|$. The third term of Hamiltonian (1) can be considered as perturbation that causes radiative transitions between spin subbands with spin flip in the presence of electromagnetic wave with vector potential $\vec{A}$.

For transition metals, atomic spin is predominantly determined by the spin of *d* electrons. Therefore, we use the formula for estimation of the exchange *sd* interaction from [34] and introduce the *s*-exchange tensor [31, 32]

$$(\vec{H}^0_{sd})_k = (\alpha_{sd})_{kl}(\vec{M}_2)_l, \quad (5)$$

where summation is performed with respect to identical indices. We use formula (3) to calculate *sd*-exchange field $\vec{H}_A$ induced by electromagnetic wave and employ formula (5) to show that vectors $\vec{H}^0_{sd}$ and $\vec{H}_A$ are collinear if the medium is isotropic. Owing to the difference of diagonal terms of the *sd*-exchange tensor (when reduced to the principal axes), fields $\vec{H}^0_{sd}$ and $\vec{H}_A$ are noncollinear and, hence, the medium is anisotropic. Only in this case, the model under study allows radiative transitions with spin flip of electrons in reduced states. This circumstance leads to noticeable changes in the analysis of the spin-injection emission in comparison with the analysis of [26].

In accordance with formulas (2)–(4), for convenience of the further analysis, we introduce vector $\vec{I}$,

which has the dimension of energy using effective *sd*-exchange field $\vec{H}_{sd}^{0}$:

$$\vec{I} = \hbar\gamma\frac{\vec{H}_{sd}^{0}}{2} = \mu_B G(\vec{p})\vec{M}_2(\vec{r}), \quad (6)$$

where $G(\vec{p}) = \alpha_{ij}(\vec{p})$ is the *sd*-exchange tensor (note that selected directions of anisotropy may emerge in thin single-crystalline films with cubic symmetry [34]). We also introduce vector $\vec{J}$ using *sd*-exchange field $\vec{H}_A$ induced by electromagnetic wave:

$$\vec{J} = \hbar\gamma\frac{\vec{H}_A}{2} = \frac{e}{2c}\left(\frac{\partial\vec{I}}{\partial\vec{\hat{p}}}\vec{A} + \vec{A}\frac{\partial\vec{I}}{\partial\vec{\hat{p}}}\right). \quad (7)$$

Vector $\vec{H}_{sd}$, which determines the *sd*-exchange field with allowance for an additional contribution of the external electromagnetic field is introduced as

$$\vec{H}_{sd}^{*} = 2(\vec{I} - \vec{J})/\hbar\gamma. \quad (8)$$

Note that the coordinate component of field $\vec{H}_{sd}$ is supplemented with time component $\vec{J}$ the oscillation frequency of which is determined by the external electromagnetic field. We assume that the frequency of the external field is equal to the transition frequency of electron with spin flip between subbands $\omega_{12} \approx (2I_z)/\hbar$. If the frequency of external field is comparable with the frequency of transition between the subbands (which is greater than the frequency of ferromagnetic resonance by several orders of magnitude), the lattice magnetization (as was mentioned) cannot follow the external field and the initial position is retained. Then, we disregard the dynamics of magnetic lattice and restrict consideration to transverse components of $\vec{A}(z,t)$, (e.g., components $A_x(z,t)$ and $A_y(z,t)$). Thus, the Hamiltonian of nonequilibrium injected electrons is determined by the exchange field that is modulated by the external radiation. The anisotropy of the exchange interaction and the presence of external field lead to a term in Hamiltonian (1) that describes the perturbation and contains off-diagonal terms that lead to spin flip.

3. ANALOG OF THE BLOCH EQUATION FOR INJECTED ELECTRONS WITH ALLOWANCE FOR ELECTRON TRANSPORT

To describe kinetics of the magnetization of injected electrons, it is expedient to consider an analog of the modified Bloch equation (with the transport term) and take into account the interaction of injected electrons with electromagnetic field. Such an equation can be used to study both relaxation and emission processes and the transport of nonequilibrium spins. Spatial nonuniformity of the spin distribution must be taken into consideration to take into account the transport term in the analog of the Bloch equation. The equation has been derived in our previous work [32].

An analog of the Bloch equation with the transport term for the magnetization of injected electrons that takes into account transport and interaction of conduction electrons with electromagnetic wave via the *sd* interaction is represented in the following way (transverse and longitudinal decays related to nonradiative spin relaxation upon interaction with medium are introduced phenomenologically):

$$\begin{aligned}\frac{d\vec{\mu}_\perp}{dt} &= \gamma(\vec{\mu}\times\vec{H}_{sd}^{*})_\perp - \nabla_k\langle\Pi_k\mu_\perp\rangle - \frac{\vec{\mu}_\perp}{\tau_\perp},\\ \frac{d\vec{\mu}_\parallel}{dt} &= \gamma(\vec{\mu}\times\vec{H}_{sd}^{*})_\parallel - \nabla_k\langle\Pi_k\mu_\parallel\rangle - \frac{\vec{\mu}_\parallel - \langle\vec{\mu}_\parallel\rangle}{\tau_\parallel},\end{aligned} \quad (9)$$

$$\Pi_k = \frac{1}{2m^*}\{\delta(\vec{r} - \vec{r}_o), \hat{p}_k\}. \quad (10)$$

Here, braces {, } denote anticommutator, $\hat{p}_k$ is the operator of the electron momentum, $\vec{\mu}_\parallel$, and $\vec{\mu}_\perp$ are the longitudinal and transverse components of the magnetic moment of the spin of injected electrons, $\tau_\parallel$ is the longitudinal (spin) relaxation time of injected electrons, $\tau_\perp$ is the relaxation time of the transverse spin component that we assume to be related to the Slonczewski–Berger mechanism [17]. The relaxation time of the transverse spin component is relatively small, and the relaxation takes place upon propagation of the wave at a distance of about several nanometers [17]. The longitudinal spin relaxation time is significantly greater, and the relaxation occurs upon wave propagation at a distance of several tens of nanometers. System of equations (9) and its solution describe the process when the longitudinal relaxation time is determined by the intensity of the processes of energy exchange between the injected electron and lattice [36], and the transverse relaxation time is determined by the rate at which the synchronous precession of magnetic dipoles is disturbed [17]. The relaxation of the transverse spin component is relatively fast on the time scale of the longitudinal relaxation. Note that the spin torque is transferred to the magnetic lattice of the ferromagnetic material upon relaxation of the transverse spin component. In this case, oscillations of magnetization or magnetization switching of the ferromagnetic material may take place [1–10] (such processes are not taken into account in this work). The further spin relaxation involves transitions of electrons between nonequilibrium quasi-Fermi levels with spin flip, which may lead to emission in terahertz and far-IR spectral ranges.

System of equations (9) shows that the state of injected nonequilibrium electrons is changed with time due to the effect of the *sd* exchange, external electromagnetic field, nonradiative relaxation and transfer of carriers, and electron transport. Below, we show that Eqs. (9) can be used to derive an equation for non-

equilibrium spin polarization in magnetic junctions with allowance for spatial inhomogeneity of the spin polarization.

4. EQUATIONS FOR NONEQUILIBRIUM SPIN POLARIZATION IN MAGNETIC JUNCTIONS

We consider the scenario in which electric current flows through the interface of homogeneous ferromagnetic materials the magnetization vectors of which (that correspond to quantization axes in the materials) make mismatch angle φ. We introduce the following notation: P_1, nonequilibrium spin polarization of the injector (ferromagnetic material from which spin-polarized electrons are injected (Fig. 1)) and P_2, equilibrium spin polarization of the working region of magnetic junction to which spin-polarized electrons are injected. We assume that the electric current flows along the x axis through the interface of ferromagnetic materials. For simplicity, we assume that electric current density j changes neither along the transverse cross section nor along the x axis of magnetic junction. Then, we assume negligibly small anisotropy of the exchange interaction (i.e., zero components of the exchange field determined by the external electromagnetic field) and the fact that the relaxation of the longitudinal component of the magnetic moment of conduction electron is fully determined by the interaction of the spin system of conduction electrons and thermostat.

The tensor of magnetization flux is defined as [34]

$$J^{ik} = \frac{i\hbar\mu_B}{2m}\sum_{\vec{p},s_1,s_2}\sigma^i_{s_1,s_2}\left(\Psi_{\vec{p},s_2}\nabla_k\Psi^*_{\vec{p},s_1} - \Psi^*_{\vec{p},s_1}\nabla_k\Psi_{\vec{p},s_2}\right) = \mathrm{Re}\sum_{\vec{p},s_1,s_2}\Psi^*_{\vec{p},s_1}\sigma^i_{s_1,s_2}\hat{v}^k\Psi_{\vec{p},s_2}, \quad (11)$$

where $\Psi_{\vec{p},s}(\vec{r},t)$ is the wave function of electron with momentum $\vec{p}$ in spin state s; μ_B is the Bohr magneton; $\vec{\sigma} = \{\sigma^x,\sigma^y,\sigma^z\}$ is the Pauli vector matrix; spin and vector indices are shown as subscripts and superscripts, respectively; and $\hat{v} = -(i\hbar/m)\nabla$ is the velocity operator. When formula for Π_k (10) and formula (11) are used, it is seen that quantities $n\langle\Pi_k\mu_\perp\rangle$ and $n\langle\Pi_k\mu_\parallel\rangle$ in formulas (11) in the above configuration represent averaged components of the magnetization flux tensor.

Let's consider transformation of the injector magnetization component in the above configuration

$$J^{zx} = \frac{i\hbar\mu_B}{2m} \times \sum_{\vec{p},s_1,s_2}\sigma^z_{s_1,s_2}\left(\Psi_{\vec{p},s_2}\frac{\partial}{\partial x}\Psi^*_{\vec{p},s_1} - \Psi^*_{\vec{p},s_1}\frac{\partial}{\partial x}\Psi_{\vec{p},s_2}\right) \quad (12)$$

upon transition to a new quantization axis. Such a transition can be due to either transition of electron to another magnetic medium with a different direction of the quantization axis or rotation (e.g., in the presence of external magnetic field) of the quantization axis of the homogeneous medium. The following matrix of transformation of spin wave functions corresponds to the rotation of coordinate axes by angle φ around the x axis [37]:

$$\hat{U}_x(\varphi) = \begin{pmatrix} \cos\frac{\varphi}{2} & i\sin\frac{\varphi}{2} \\ i\sin\frac{\varphi}{2} & \cos\frac{\varphi}{2} \end{pmatrix}. \quad (13)$$

Hence, the component of magnetization tensor is transformed in the following way upon transition to a new quantization axis:

$$J^{z'x} = J^{zx}\cos\varphi, \quad (14)$$

due to the condition $J^{yx} = 0$. In this case, we also obtain orthogonal component $J^{y'x}$, given by

$$J^{y'x} = J^{zx}\sin\varphi. \quad (15)$$

Thus, the magnetization flux (and the corresponding spin flux) from injector J^{zx} passes to the medium with a new quantization axis and is transformed into a flux that has component $n\langle\Pi_k\mu_\parallel\rangle = J^{zx}\cos\varphi$ polarized along the new quantization axis and component $n\langle\Pi_k\mu_\perp\rangle = J^{zx}\sin\varphi$ polarized orthogonally to the quantization axis. Recall that significantly different spin relaxation times correspond to the longitudinal and transverse polarizations. Hence, only longitudinal polarization survives outside a layer with a thickness on the order of the length of the transverse relaxation (Slonczewski–Berger layer). Below, we consider a region that lies outside such a layer, so that the term with vector product $\vec{\mu}\times\vec{H}^*_{sd}$ is eliminated in Eqs. (9). Thus, to derive an equation for nonequilibrium spin polarization, we employ only the second equation of system (9).

Magnetic state of conduction electrons is described using local spin polarization

$$P(x) = \frac{\vec{\mu}_\parallel(x)}{\mu_B} = \frac{n_+(x) - n_-(x)}{n}, \quad (16)$$

where $n_{+(-)}(x)$ are the partial concentrations of conduction electrons with spins that are parallel (antiparallel) to the quantization axis of the corresponding ferromagnetic material and $n = n_+(x) + n_-(x)$ is the total concentration that is assumed to be independent of coordinate owing to neutrality of metal. Taking into account expression (12), we introduce spin flux density and obtain

$$J_s(x) = J^{zx}/n = \langle\Pi_k\mu_\parallel\rangle = \frac{\mu_B}{en}[j_+(x) - j_-(x)], \quad (17)$$

where $j_\pm$ are the partial densities of electric (charge) current.

With allowance for expressions (16) and (17), the second equation of system (9) is represented under stationary conditions ($\partial\bar{\mu}_{\parallel}/\partial t = 0$) as

$$\nabla J_s(x) = -\mu_B \frac{P(x) - \bar{P}}{\tau}. \quad (18)$$

In the presence of electric field $E(x)$ and gradients of partial concentrations, we obtain partial densities of electric (charge) current

$$j_{\pm} = e\mu_{\pm} n_{\pm}(x)E(x) - eD_{\pm}\frac{dn_{\pm}}{dx}, \quad (19)$$

where $\mu_{\pm}$ and $D_{\pm}$ are the partial mobilities and diffusion coefficients of electrons, respectively.

Total density $j = j_+(x) + j_-(x)$ in the stationary state is independent of x. Representing quantity $E(x)$ in terms of j with allowance for expression (17), we obtain the spin flux density

$$J_s(x) = (\mu_B/en)\{Q(P)\, j - enD(P)\nabla P\}, \quad (20)$$

where

$$Q(P) = \frac{\mu_+ - \mu_- + (\mu_+ + \mu_-)P}{\mu_+ + \mu_- + (\mu_+ - \mu_-)P},$$
$$D(P) = \frac{\mu_+ D_- + \mu_- D_+ + (\mu_+ D_- - \mu_- D_+)P}{\mu_+ + \mu_- + (\mu_+ - \mu_-)P}. \quad (21)$$

We assume that carriers have identical mobilities and diffusion coefficients in the subbands: $\mu_- = \mu_+ = \mu$, $D_- = D_+ = D$. Then, we have $Q(P) = P$, $D(P)$ D, and substitution of expression (20) in expression (18) with allowance for conservation of electric charge $\nabla\vec{j} = 0$ yields the following equation [26]:

$$\frac{d^2P}{dx^2} - \frac{j}{j_D l}\frac{dP}{dx} - \frac{P - P_2}{l^2} = 0, \quad (22)$$

where τ is the spin relaxation time, $j_D = enD/l = enl/\tau$ is the diffusion current density of electrons, $l = \sqrt{D\tau}$ is the spin relaxation length, and n is the electron concentration in metal.

5. NONEQUILIBRIUM SPIN POLARIZATION IN MAGNETIC JUNCTIONS

The solution to Eq. (22) is written as [29]

$$P(x) = C_1 \exp(\lambda_1 x) + C_2 \exp(\lambda_2 x) + P_2, \quad (23)$$

where

$$\lambda_{1,2} = \frac{j}{2j_D l} \mp \sqrt{\left(\frac{j}{2j_D l}\right)^2 + \frac{1}{l^2}}. \quad (24)$$

Here, subscripts 1 and 2 correspond to signs "–" and "+", respectively. The solution to Eq. (22) can be simplified under the condition:

$$j/j_D \ll 1. \quad (25)$$

For a spin diffusion length of $l = \sqrt{D\tau} \sim 3\times10^{-6}$ cm, a concentration of conduction electrons in metal of $n \sim 10^{22}$ cm^{-3}, and a longitudinal (spin) relaxation time of injected electrons of $\tau_{\parallel} \sim 3\times10^{-13}$ s [6, 7], we estimate electron diffusion current density as $j_D = enl/\tau \sim 1.6\times10^{10}$ A/cm^2. Normally, the maximum current densities in magnetic junctions in experiments are less than the estimated current density by an order of magnitude. Therefore, condition (25) is reliably satisfied. Then, solution (23) is represented as

$$P(x) = C_1 \exp(-x/l) + C_2 \exp(x/l) + P_2. \quad (26)$$

We have $C_2 = 0$, since $P(\infty) = P_2$. Integration constant C_1 is determined from the condition for continuity of the spin flux at the interface of the two ferromagnetic materials [6]

$$J_s(0) = \frac{\hbar}{2e}\left\{jP(0) - j_D l \left.\frac{dP}{dx}\right|_{x=0}\right\} = \frac{\hbar}{2e} P_1 j\cos\varphi. \quad (27)$$

Boundary condition (27) is valid under the condition that the spin polarization of injector remains unchanged or is weakly perturbed, which is possible for appropriately chosen materials [6]. In addition, we assume that the effect of electron flux reflected from the interface on the spin polarization of electrons having passed to the working region of the magnetic junction can be disregarded.

Substituting expression (23) in expression (27), we obtain

$$C_1 = \frac{P_1\cos\varphi - P_2}{j - j_D l\lambda_1} j. \quad (28)$$

With allowance for formula (28), the distribution of spin polarization of conduction electrons in the working region is written as

$$P(x) = P_2 + \frac{P_1\cos\varphi - P_2}{j - j_D\lambda_1 l} j\exp(\lambda_1 x). \quad (29a)$$

Formula (29a) is valid at $j/j_D \ll 1$, and can be simplified as

$$P(x) = P_2 + \frac{P_1\cos\varphi - P_2}{j + j_D} j\exp(-x/l). \quad (29b)$$

In the opposite case $j \to \infty$ ($j \gg j_D$), formulas (24) and (29a) yield

$$P(x) = P_2 + \frac{P_1\cos\varphi - P_2}{j + (j_D/2j)^2} j\exp(-j_D x/4jl). \quad (29c)$$

For $j/j_D \to 0$, formula (29c) shows that the spin polarization of injected electrons becomes indepen-

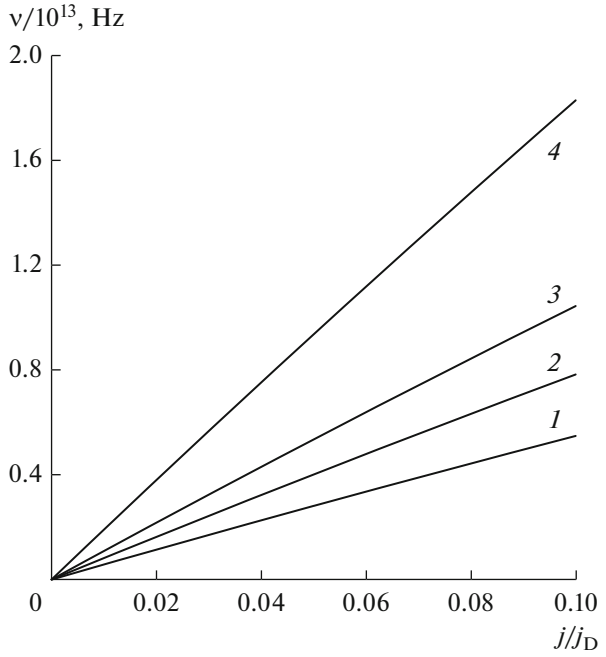

Fig. 3. Plot of frequency of spin-injection radiation ν vs. relative current density j/j_D at an angle of $\varphi = 120°$ between magnetizations of the first and second layers of magnetic junction and polarization-degree ratios of the first and second layers of P_1/P_2 = (*1*) 0.1, (*2*) 1, (*3*) 2, and (*4*) 5.

dent of distance from the injector, and we have $P(x) \to P_1 \cos\varphi$. Thus, the spin polarization in the working region of the magnetic junction cannot be greater than the spin polarization in the injected flux.

Formulas (29a)–(29c) show that, as was expected, the spin polarization monotonically approaches the equilibrium level for the working medium with an increase in the distance from the injector. The maximum deviation of the spin polarization from the nonequilibrium polarization $\Delta P = P - P_2$ is reached at the interface:

$$\Delta P(0) = \frac{P_1 \cos\varphi - P_2}{j - j_D l \lambda_1} j. \qquad (30a)$$

For $j/j_D \ll 1$, formula (30a) is written as

$$\Delta P(0) = \frac{P_1 \cos\varphi - P_2}{j + j_D} j. \qquad (30b)$$

6. SPIN-FLIP ENERGY TRANSITIONS OF ELECTRONS

Nonequilibrium electrons have nonequilibrium quasi-Fermi levels ε_{F+} and ε_{F-} that are determined relative to the bottom of the lower spin subband (Fig. 2). Under nonequilibrium conditions, electron transitions are possible from occupied states that are lower than quasi-Fermi level ε_{F+} (e.g., for the minor spin subband) to free electron states that are higher than quasi-Fermi level ε_{F-} for the major spin subband (Fig. 2). If quasi-Fermi level ε_{F+} is shifted relative to the equilibrium position of the Fermi level in the working region by $\Delta\varepsilon_{F+} = \varepsilon_{F+} - (\varepsilon_{F0+} - eV/2)$, and quasi-Fermi level ε_{F-} is shifted by $\Delta\varepsilon_{F-} = (\varepsilon_{F0-} - eV/2) - \varepsilon_{F-}$, the maximum radiation frequency can be represented as a sum of quantities $\Delta\varepsilon_{F+}$ and $\Delta\varepsilon_{F-}$, divided by the Planck constant [26]:

$$\nu \equiv \frac{\Delta\varepsilon_{F+} + \Delta\varepsilon_{F-}}{2\pi\hbar} = \frac{|\varepsilon_{F+} - \varepsilon_F| + |\varepsilon_{F-} - \varepsilon_F|}{2\pi\hbar}. \qquad (31)$$

Here, we use modulus, to make formula valid for any sign of quantity ΔP.

Quasi-levels of the subbands depend on the nonequilibrium spin polarization and exchange energy splitting (I_2):

$$\varepsilon_{F+} = \frac{\hbar^2}{2m}\left(3\pi^2 n\right)^{2/3}\left[\left(\frac{1 - P_2 + |\Delta P|}{2}\right)^{2/3}\right] + 2I_2,$$

$$\varepsilon_{F-} = \frac{\hbar^2}{2m}\left(3\pi^2 n\right)^{2/3}\left[\left(\frac{1 + P_2 - |\Delta P|}{2}\right)^{2/3}\right].$$

The formulas are written with allowance for the fact that a phase volume of $(2\pi\hbar)^3$, corresponds to one electron and the number of electrons with the same direction of spin $n_{-,+} = n(1 \pm P)/2$ is equal to the volume of a sphere with a radius that is equal to the Fermi momentum divided by $(2\pi\hbar)^3$. For convenience, we represent variations in the spin polarization of quasi-levels in terms of equilibrium value (P_2) and additional nonequilibrium term (ΔP):

$$\Delta\varepsilon_{F+} = \frac{\hbar^2}{2m}\left(3\pi^2 n\right)^{2/3}\left|\left(\frac{1 - P_2 - \Delta P}{2}\right)^{2/3} - \left(\frac{1 - P_2}{2}\right)^{2/3}\right|, \qquad (32)$$

$$\Delta\varepsilon_{F-} = \frac{\hbar^2}{2m}\left(3\pi^2 n\right)^{2/3}\left|\left(\frac{1 + P_2}{2}\right)^{2/3} - \left(\frac{1 + P_2 + \Delta P}{2}\right)^{2/3}\right|. \qquad (33)$$

In formulas (32) and (33), the sign of variations in quasi-levels depends on the sign of nonequilibrium spin term, so that the variations can be positive and negative.

Figure 3 presents the dependence of radiation frequency ν on reduced current density calculated with the aid of formulas (31)–(33). It is seen that a frequency interval of the spin-injection radiation of 5–20 THz corresponds to the current density that is greater than the diffusion current density by approximately two orders of magnitude. Using estimated diffusion current density of $j_D \sim 10^{10}$–10^{11} A/cm^2, we obtain a current density in the magnetic junction of $j \sim 10^8$–10^9 A/cm^2 that can be reached in experiments in the absence of

the thermal breakdown of magnetic contact. Lower radiation frequencies correspond to lower current densities. We must also take into account decay of electromagnetic wave that is predominantly determined by the absorption of waves by free carriers in materials of magnetic junctions. The results of [24] show that a threshold current density of $j_{pr} \sim 10^7$–10^8 A/cm^2 must be surpassed to provide a radiation power level that corresponds to loss related to damping and thermal noise.

7. PROBABILITY OF QUANTUM TRANSITIONS

Equations (9) can be used to find the probability of transitions (number of quantum transitions per unit time) [31]:

$$W \approx n_{en}\left|\left(P_z(0) - \frac{P_z^e}{(1+\gamma^2 H_A^{2(\perp)}\tau_\perp\tau_\parallel)}\right)\left(\frac{1+\gamma^2 H_A^{2(\perp)}\tau_\perp\tau_\parallel}{\tau_\parallel}\right)\right|. \quad (34)$$

Here, n_{en} is the concentration of injected spin-nonequilibrium electrons in metal, $H_A^{(\perp)} = (H_{sd}^{2(x)} + H_{sd}^{2(y)})^{1/2}$ is the modulus of the transverse component of exchange field, $P_z(0)$ is the nonequilibrium spin polarization at the initial moment, and P_z^e is the equilibrium spin polarization in ferromagnetic material.

At a relatively low Q factor, when $\omega_{pr}\tau_{eff} \ll 1$ ($\tau_{eff} = \sqrt{\tau_\perp\tau_\parallel}$, $\omega_{pr} \approx \gamma H_A^{(\perp)}$ is the precession frequency), the probability of quantum transitions is given by formula (34):

$$W \approx n_{en}\left|\frac{P_z(0) - P_z^e}{\tau_\parallel}\right|. \quad (35)$$

For a high Q factor, when $\omega_{pr}\tau_{eff} \gg 1$, approximate formula (34) coincides with the numerical solution to system (9) at $t \approx 10\tau_\parallel$. In this case, the probability of transitions given by formula (34)

$$W \approx n_{en}\left|10\left(P_z(0)\right)\gamma^2 H_A^{2(\perp)}\tau_\perp\right| \approx 10^{37} \text{ s}^{-1}\text{ cm}^{-3} \quad (36)$$

is determined by transverse time of spin relaxation.

Note several properties of the relaxation processes that follow from formula (34). When $\omega_{pr}\tau_{eff} \gg 1$ and $t \gg \tau_\parallel/\omega_{pr}\tau_{eff}$, the population of the upper and lower electronic sublevels become equal. At a high-Q precession (i.e., in the presence of relatively strong external electromagnetic field $H_A^{(\perp)}/H_{sd}^{*(z)} = 0.1$), the rate of energy transfer from the field may be higher than the rate at which the oscillation energy of magnetic moment is transferred to the thermostat. Thus, the energy exchange between the electromagnetic field and oscillator is balanced and the system reaches an equilibrium state with respect to the external field, so that the saturation takes place [38].

The probability of spin-flip transitions of conduction electrons upon interaction with thermostat is given by formula (35) (below, we determine the contribution of radiative transitions). On the assumption that the quantum yield equals one (upper-bound estimation), formula (35) yields a number of radiative transitions of $W_r \approx 10^{31}$ s^{-1} cm^{-3} for a density of nonequilibrium electrons of $n_{en} \approx 10^{19}$ cm^{-3} (a current density of 10^8 A/cm^2, see formula (29)). In this case, liberated radiation power per unit volume P_i is determined as the number of transitions multiplied by photon energy $\hbar\omega$. For ω_{12}= 30 THz, we have

$$P_i = \hbar\omega W_r \approx 10^{11} \text{ W cm}^{-3}. \quad (37)$$

The volume of the active region of the emitter is determined by the area of magnetic junction and spin relaxation length. For a junction diameter of about 1 μm and a spin relaxation length of 10 nm, we obtain a volume of the active region of $V = \pi d^2 l \approx 10^{-15}$ cm^3. Thus, the maximum radiation power that can be emitted in the active region is $P_i \approx 10^{-4}$ W in accordance with formula (37). Such a high estimated power is obtained with disregard of the fact that, in practice, the quantum yield is significantly less than one and the number of nonequilibrium electrons is not equal to a maximum number of 10^{19} cm^{-3} at the given current density [29].

CONCLUSIONS

We have derived an equation for nonequilibrium spin polarization in magnetic junctions based on the equation for dynamics of magnetic moment averaged over the ensemble of nonequilibrium spin-injected electrons with allowance for spatial nonuniformity of the distribution. Using the solution of the equation for nonequilibrium spin polarization at the interface of two ferromagnetic materials with the current flow in the contact region, we have obtained formulas for calculation of photon emission and absorption in the case of indirect energy transitions of electrons with spin flip. It was shown that at a current density that exceeds a threshold level, radiation can be emitted from the contact region and the corresponding frequencies belong to the terahertz range. In general, the frequencies of photon emission and absorption for indirect

energy transitions of electrons with spin flip are determined by the current density, the angle between the magnetization directions of the injector and working region, and spin polarization of conduction electrons. At current densities that substantially exceed the diffusion current density, the maximum possible nonequilibrium spin polarization is obtained in the working layer and the corresponding frequencies of photon emission and absorption reach saturation. The results can be used for development of compact sources of terahertz radiation.

FUNDING

This work was supported in part by the Russian Foundation for Basic Research (project nos. 18-57-76001 ERA_a, 18-29-27020 mk, and 19-29-03015 mk) and State Program no. 075-00475-19-00.

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Translated by A. Chikishev