Non-equilibrium spin waves in paramagnetic metals

A. A. Zyuzin and A. Yu. Zyuzin

A.F. Ioffe Physico-Technical Institute of Russian Academy of Sciences - 194021 St. Petersburg, Russia

received 11 March 2010; accepted in final form 14 June 2010
published online 15 July 2010

PACS 75.30.Ds – Spin waves
PACS 71.27.+a – Strongly correlated electron systems; heavy fermions
PACS 72.25.Hg – Electrical injection of spin-polarized carriers

Abstract – We theoretically study the effect of exchange interaction on the non-equilibrium spin waves in disordered paramagnetic metals under the spin injection condition. We show that the gapless spectrum of spin waves, describing the spin precession in the absence of the applied magnetic field, changes sign to negative on the paramagnetic side near the ferromagnet-paramagnet phase transition. The damping of spin waves is small in the limit when electron-electron exchange energy is larger than the inverse electron mean free time, while in the opposite limit the propagation of spin waves is strongly suppressed. We discuss the amplification of the electromagnetic field by the non-equilibrium spin waves.

Introduction. – Effects related to non-equilibrium spin accumulation in metals and semiconductors attract considerable interest in the field of spin phenomena (see for a review [1,2]). The spin accumulation in nonmagnetic materials can be produced by different types of pumping such as, for example, optical orientation and electrical methods. In the optical spin orientation method [3] the electron’s orbital momentum is oriented by the absorption of circularly polarized electromagnetic field and as a result spin becomes polarized through the spin-orbit interactions. The electrical spin injection method is based on transfer of spin-polarized electrons from the ferromagnet into the paramagnet. The mechanism of spin injection from ferromagnet into semiconductors and metals was first proposed by Aronov and Pikus [4,5]. It was shown that the electrical current through the contact between a ferromagnet and a paramagnet will produce the non-equilibrium magnetization in the paramagnetic region on the scale of the spin-diffusion length. The non-equilibrium spin magnetization was experimentally studied in metals by Johnson and Silsbee [6,7] by voltage measurement.

Despite of extensive study of various aspects of spin dynamics in paramagnets [1,2], little attention was paid to the effect of electron-electron exchange interaction on the non-equilibrium spin transport. Aronov [8] predicted the existence of the long-lived spin-waves excitations in the clean paramagnets with the non-equilibrium spin polarization under spin injection. The spectrum of these spin waves has a quadratic power low dependence on the wave vector $\propto q^2$ at small $q$, while the Landau damping of spin waves is small compared to theirs frequency. Interestingly, spin waves in non-equilibrium system have negative frequency, which means that the excitation of these spin waves lowers the energy of the system. Motivated by this research, we extend the study to the case of strongly enhanced and disordered paramagnetic metals with the non-equilibrium spin polarization.

Berger [9] and Slonczewski [10] calculated the spin-transfer torque effect in thin-film ferromagnetic-semimagnetic metal multilayer system. They predicted the stimulated emission of spin waves in ferromagnetic layers due to negative Gilbert damping at high values of current density applied to the multilayer. M. Tsoi et al. [11] first observed such current-induced excitations in a magnetic multilayer in magnetoresistance measurements. Recently, the direct measurements of both the magnitude and the direction of the spin-transfer torque in a magnetic tunnel junction were given in papers (see [12,13] and references therein). We argue that the effects related to the non-equilibrium spin excitations in the paramagnetic layers also might become interesting.

In the present paper, we show that in contrast to the stimulated emission of spin waves in the ferromagnetic layers of the magnetic multilayer systems [9,10], the non-equilibrium transverse spin waves in paramagnetic metals are stable under spin injection. On the other hand the excitation of these spin waves decreases the energy of the system. Indeed, we found that there is a crossover from positive to negative sign of the frequency in the vicinity of the ferromagnet-paramagnet phase transition.
The real part of the transverse spin-waves frequency in the paramagnetic metal has negative sign, while the imaginary part describing the damping of spin waves is negative. We found that in the regime when the exchange energy of electrons is smaller than the inverse mean free time the propagation of spin waves is strongly suppressed. Interestingly, the electromagnetic field might be amplified via coupling to the non-equilibrium spin waves. Clearly, one deals with the system driven by the external spin injection source to the highly excited spin correlated state. The energy might be extracted from this state via creation of the spin waves. Based on these results, we propose an electromagnetic field amplification in the spin-wave resonance experiment.

Definitions. – Let us consider the tunneling contact between a ferromagnetic metal with spontaneous magnetization and a paramagnetic metal. The current flow through the contact produces the non-equilibrium magnetization in the paramagnetic region. The realization of spin injection is well studied experimentally. In addition, we take into account the external magnetic field applied to the paramagnet leading to the Zeeman effect. For simplicity, we assume the direction of magnetic field is parallel to the non-equilibrium magnetization in the paramagnetic metal. We suggest the size of the paramagnet is much smaller than the typical spin relaxation length corresponding, for example, to the spin-orbit scattering. Therefore, the electron interactions with impurities are assumed to be spin independent.

In order to obtain the spectrum of transverse spin waves in the paramagnet we will calculate the poles of the magnetic susceptibility using Keldysh formalism. The equation for the Keldysh function $G_K$ and the retarded and advanced Green’s functions $G_R$, $G_A$ is presented in fig. 1, left.

The retarded and advanced Green’s functions in the Fourier representation are given by

$$ G_{R,A}^{\uparrow,\downarrow}(p,\omega) = [\omega - E_p^{\uparrow,\downarrow} + \mu \pm i/2\tau]^{-1} $$

where $\tau$ is the mean free time and $\mu$ is the chemical potential. The energy of spin-up and spin-down electrons is given as

$$ E_p^{\uparrow,\downarrow}(p) = \frac{p^2}{2m} + \omega_z - \frac{i}{2} \lambda \int \frac{dk d\omega}{(2\pi)^3} G_{R}^{\uparrow,\downarrow}(k, \omega), $$

where $\lambda$ is the electron-electron exchange coupling constant. Last term in expression (2) describes contribution of the short-range electron-electron exchange interactions to the spin splitting, fig. 1, left, see [14]. The “Hartree term” has negative sign and consists of two contributions corresponding to two spin-up and spin-down electron states at the bubble. The “Fock term” to the self-energy of electrons is positive and has one spin contribution which cancels the Hartree contribution with equal spins. Correspondingly, there is only one term left, shown in fig. 1, left, with different spins denoted as $\alpha$ and $\beta$.

We assume that the free electron dispersion is parabolic $p^2/2m$ with an effective mass $m$. We also introduce here an external magnetic field which contributes the Zeeman energy $\omega_z = \mu_B H_z$ to definition (2). We suggest that the magnetic field is applied in the $z$-direction and parallel to the non-equilibrium magnetization in the paramagnet.

Spin-polarized current flow through the ferromagnet-paramagnet interface leads to the different densities of spin-up and spin-down electrons in the paramagnetic metal [5] on the scale of spin relaxation length. Since we assume the spin relaxation processes to be weak, we suggest the system under consideration is in the energy equilibrium state while having non-equilibrium spin accumulation. The resulting spin accumulation is uniform over the paramagnetic sample and is proportional to the value of spin-polarized current. The electron chemical potential shifts for spin-up and spin-down electrons $\mu^{\uparrow,\downarrow} = \pm \delta \mu/2$.

Then the Keldysh function does not depend on the spatial coordinates of the system and one has

$$ G_{R}^{\uparrow,\downarrow}(p, \omega) = (1 - 2f^{\uparrow,\downarrow}(\omega))[G_{R}^{\uparrow,\downarrow}(p, \omega) - G_{A}^{\uparrow,\downarrow}(p, \omega)] $$

where Fermi distribution function at a given temperature $T$ is $f^{\uparrow,\downarrow}(\omega) = (e^{(\omega - \mu^{\uparrow,\downarrow})/T} + 1)^{-1}$. Due to the electron exchange interaction the energy of electrons depends on the non-equilibrium chemical potential shifts.

The ladder approximation [14-16] for the transverse dynamical magnetic susceptibility is shown in fig. 1, right, where we assume short-range electron-electron interactions. The spectrum of non-equilibrium spin waves is determined by the poles of the susceptibility $\chi^{+-}(q, \Omega)$ for the circular spin component $s^{\uparrow,\downarrow} = s_\alpha + is_\beta$.

The dynamical magnetic susceptibility can be evaluated in the leading order in $1/\mu \tau$ by summing the ladder vertex correction to the polarization bubble due to exchange interactions between electrons and impurity scattering. The general result for the dynamical susceptibility takes the form

$$ \chi^{+-}(q, \Omega) = 2\mu^2 \nu \frac{\Pi(q, \Omega)}{1 - \lambda q \Pi(q, \Omega)} $$

where $\nu$ is the inelastic mean free time of impurities and $\Pi(q, \Omega)$ is the inelastic electron-impurity Green’s function.
where \( \nu \) is the electron density of states per spin direction and \( \Pi(q, \Omega) \) is determined through the non-equilibrium Green’s functions as

\[
\Pi(q, \Omega) = \frac{i}{2\nu} \int \frac{dp dq d\omega}{(2\pi)^3} \left[ \frac{G^r_K(p+, \omega_+)}{D_{AA} D_{RR}} + \frac{G^r_A(p-, \omega_-) G^r_K(p+, \omega_+)}{D_{AR} D_{AA}} \right],
\]

(5)

where subscripts \( p_\pm = (p \pm q/2)^2/2m \) and \( \omega_\pm = \omega \pm \Omega/2 \) define the electrons momentum and frequency, while functions \( D_{XY} \) are given as

\[
D_{XY} = \frac{1}{2\nu \tau} \int \frac{dp dq d\omega}{(2\pi)^3} G^r_K(p-, \omega_-) G^r_Y(p+, \omega_+).
\]

Evaluating expression (4), we suggest the exchange coupling \( \lambda(p) \) to be constant. Further we will discuss the dependence of \( \lambda(p) \) on momentum and obtain the corresponding spin wave’s dispersion.

The spin polarization \( \mathcal{P} \) in the paramagnetic region is defined as

\[
\mathcal{P} = -\frac{i}{2N} \int \frac{dp dq d\omega}{(2\pi)^3} \left[ G^r_Y(p, \omega) - G^r_K(p, \omega) \right],
\]

(7)

where \( N \) is the total electron concentration. Equation (7) is equivalent to the condition \( 1 - \lambda \Pi(0, \Omega = 2\omega_z) = 0 \). In equilibrium case at \( \omega_z = 0 \) it determines the paramagnetic-ferromagnet transition point.

In general, polarization \( \mathcal{P} \) consists of equilibrium part induced by Zeeman effect and non-equilibrium part appearing due to \( \delta \mu \) shift. The non-equilibrium part of polarization can be either positive or negative depending on the direction of the injection current. It is convenient to introduce the energy corresponding to \( \mathcal{P} \)

\[
\Omega_{ex} = \lambda N \mathcal{P}.
\]

(8)

The value \( |\Omega_{ex}| \ll \mu \) is the exchange energy of electrons. Given the Zeeman splitting \( 2\omega_z \) and the shift \( \delta \mu \) of quasi-chemical potential, one obtains from (3) and (7) the equation for polarization, or equivalently for \( \Omega_{ex} \)

\[
\Omega_{ex} = \lambda \nu \Omega_{ex} + \delta \mu + 2\omega_z \left[ 1 - \frac{1}{6} \left( \frac{\Omega_{ex} + \delta \mu + 2\omega_z}{2\mu} \right)^2 \right].
\]

(9)

Note that we used the particle conservation condition under the spin injection and the parabolic electron dispersion in deriving eq. (9).

The spontaneous polarization vanishes at \( \lambda \nu < 1 \) and eq. (9) shows that the sign of \( \Omega_{ex} \) coincides with the sign of \( \delta \mu + 2\omega_z \) and in the \( \lambda \nu \ll 1 \) limit one obtains

\[
\Omega_{ex} = \frac{\lambda \nu}{1 - \lambda \nu} (\delta \mu + 2\omega_z).
\]

(10)

At \( 1 - \lambda \nu \ll 1 \) the exchange energy is much larger than \( |\delta \mu + 2\omega_z| \) and from eq. (9) it follows that:

\[
\delta \mu + 2\omega_z \approx \Omega_{ex} \frac{1}{6} \left( \frac{\Omega_{ex}}{2\mu} \right)^2.
\]

(11)

If the Stoner criterion of ferromagnetism \( \lambda \nu > 1 \) is realized then eq. (9) has the nonzero solution even in the absence of the external field and pumping [15]

\[
\frac{1}{6} \left( \frac{\Omega_{ex}}{2\mu} \right)^2 = 1 - \frac{1}{\lambda \nu}
\]

(12)

**Gapless spin waves.** – Now let us discuss the spin-waves spectrum dependence in the regimes of both spontaneous and non-equilibrium polarizations. The spectrum of spin-wave excitations is defined by the poles of the dynamical magnetic susceptibility \( \chi^{+-}(q, \Omega) \) in eq. (4), namely

\[
1 - \lambda \nu \Pi(q, \Omega) = 0.
\]

We obtain the expression for the real part of spectrum in the limit \( |\Omega_{ex}| \tau > 1 \) at small momentum and frequency \( |\Omega_{ex}| >> q v_F, |\Omega| \), where \( v_F \) is the Fermi velocity, the expression

\[
\Re \Omega = 2\omega_z \sign(n \mathbf{H}) = \frac{\lambda \nu}{3} \left( \frac{q v_F}{\Omega_{ex}} \right)^2 \left[ \delta \mu + 2\omega_z \right]
\]

\[
- \left( \Omega_{ex} - (\omega_z + \delta \mu/2) \right) \left( \frac{\Omega_{ex} + \delta \mu + 2\omega_z}{4 \mu} \right)^2,
\]

(13)

where \( n \) is the unit vector in the direction of the magnetization in the paramagnetic metal. This expression should be supplemented with eq. (9) in order to obtain the spectrum for the different regimes of pumping and values of the exchange interaction \( \lambda \nu \).

The Zeeman term in eq. (13) has negative sign so long as the direction of magnetization in the paramagnetic metal \( n \) is opposite to the direction of the magnetic field \( \mathbf{H} \). The spectrum of spin waves becomes gapless in the limit \( \omega_z = 0 \). Under this circumstance the polarization in the paramagnet is defined by the spin injection, \( \delta \mu \), only.

**Spin-waves frequency sign change.** – In this section, we consider the case where \( \omega_z = 0 \). Firstly, assume that the exchange coupling \( \lambda \nu > 1 \) and the system is in the ferromagnetic state. Using eq. (9) one sees that \( \Omega_{ex} \gg \delta \mu \) and the nonlinear term in the square brackets of eq. (13) is predominant. Taking into account the dissipation one obtains the magnon spectrum in the ferromagnet [15]

\[
\Omega = \frac{\lambda \nu}{3} \left( \frac{q v_F}{4 \mu} \right)^2 \left( |\Omega_{ex}| - \frac{\lambda \nu D q^2}{|\Omega_{ex}| \tau - i} \right),
\]

(14)

where \( D = v_F^2 \tau/3 \) is the diffusion coefficient. The real part of the spectrum is given by Goldstone mode, while the damping of spin wave is very small and proportional to the quartic term of the wave vector [17].

Now let us turn to the paramagnetic side, \( \lambda \nu < 1 \), of the ferromagnet-paramagnet phase transition when the spin polarization is induced due to spin injection. In the vicinity of phase transition point

\[
1 - \lambda \nu \ll 1
\]

(15)

Non-equilibrium spin waves in paramagnetic metals

67007-p3
the sign of the wave-vector–dependent part of the spectrum, eq. (13), stays positive as it is in the ferromagnetic phase. However, we find the sign change crossover of the $q$-dependent part of the spectrum below some critical value of $\lambda \nu$ (see also fig. 2). Indeed, in this regime eqs. (9), (11) reduce to

$$\delta \mu \approx \Omega_{ex} \frac{1}{6} \left( \frac{|\delta \mu|}{4\mu_0} \right)^2.$$  \hspace{1cm} (16)

Combining eq. (16) with eq. (13) one sees that when parameter $\lambda \nu \approx 1$ smaller is than the value

$$\lambda \nu \approx 1 - \frac{1}{3} \left( \frac{|\delta \mu|}{4\mu_0} \right)^{2/3}$$ \hspace{1cm} (17)

the wave-vector–dependent part of the spectrum in eq. (13) changes sign and becomes negative. More generally, the width of the positive-frequencies region on the paramagnetic side of the phase transition strongly depends on the details of the free electron dispersion, which was assumed parabolic.

Finally, in the limit of small $\lambda \nu \ll 1$ when the exchange energy of electrons is small compared to $\delta \mu$ and $2\omega_z$, we recover the result of Aronov [8]

$$\Omega \approx -\frac{(qv_F)^2}{3|\Omega_{ex}|}.$$ \hspace{1cm} (18)

Here $|\Omega_{ex}|$ is defined by eq. (10) at $\omega_z = 0$. This dependence is presented in fig. 3. Note that the frequency is inversely proportional to the exchange coupling and has the opposite sign compared to the case of equilibrium ferromagnet, eq. (14).

Accounting for electron scattering by impurities and Zeeman energy gives expression in limit $\lambda \nu \ll 1$

$$\Omega = 2\omega_z \text{sign}(n\mathbf{H}) + \frac{1 - \lambda \nu}{|\Omega_{ex}|} D q^2.$$ \hspace{1cm} (19)

Here the value of the exchange energy $|\Omega_{ex}|$ is given by eq. (10). Here we do not impose limitations on $|\Omega_{ex}| \tau$. The dispersion eq. (19) in the limit $\Omega_{ex} \tau > 1$ is shown in fig. 3. The inset illustrates regimes when the magnetization in the paramagnetic metal is in the same sign(nH) > 0 or in the opposite direction sign(nH) < 0 to the magnetic field. The damping of spin-wave excitations is defined by the imaginary part of the frequency $\Omega$, which in the limit $|\Omega_{ex}| \tau > 1$ is much smaller than the real part of $\Omega$. The damping of the spin-wave excitations increases in the limit $|\Omega_{ex}| \tau < 1$ and is defined by the diffusion coefficient modified by exchange interactions by a factor of $1 - \lambda \nu$.

In the case of the equilibrium regime in the absence of the spin injection, $\delta \mu = 0$, eq. (19) coincides with the results of [18,19]. Note that in the equilibrium regime the first Zeeman term in eq. (19) is positive and, as a result, the real part of the total frequency is also positive. Actually, the gapped Zeeman mode is related to the total spin precession around the external magnetic field $\mathbf{H}$. Note that the direction of the total magnetization in the paramagnet depends both on the non-equilibrium magnetization produced by spin injection and on the equilibrium part due to external magnetic field. Naturally, the total magnetization can be either in the direction of the applied magnetic field or in the opposite direction.

Let us estimate the typical values of the wave vector $q$ and the frequency $|\Omega|$ of the spin waves. In order to investigate the properties of the non-equilibrium spin waves, the typical length of the magnetic system $L = 1/q$ has to be smaller than the value of the spin-diffusion length $\ell_s$. Otherwise, the propagation of the spin waves in the paramagnet will be strongly suppressed by the spin-flip scattering.

Due to electrical current limitation it is difficult to achieve high enough non-equilibrium spin polarization in metallic systems. The possible candidate to observe the non-equilibrium spin waves might be a semiconductor structure, where polarization of order $\Omega_{ex}/\mu \sim 1$ at $\Omega_{ex}$ in meV range can be created ([20] and references therein). At Fermi velocity $v \sim 10^6$ cm/s maximum $q < 10^6$ cm$^{-1}$ might
be much larger than the inverse spin-diffusion length $\ell_s \sim 1 \mu m$, which is a reasonable value for the semiconductor structure. At $q \sim 10^6 cm^{-1}$ we estimate $\Omega$ in GHz region.

**Gapped modes.** – Now let us discuss the dependence of the exchange coupling $\lambda(p - k)$ on the momentum. Setting $p$ and $k$ to the Fermi momentum $p_F$ one can expand the interaction function over Legendre polynomials \[14\] as

$$\lambda(p - k) = \sum_{\ell \geq 0} \lambda_\ell P_\ell(\cos \theta),$$

(20)

where $\theta$ is the angle between the vectors $p$ and $k$. In the zero wave vector case, $q = 0$, the transverse magnetic susceptibility diverges if the following condition is satisfied

$$1 - \lambda_\ell \Omega_\ell(q = 0, \Omega_\ell) = 0.$$  

(21)

As a result, the spectrum of the spin waves is given by the series of gapped modes both in ferromagnet and paramagnet regimes

$$\Omega_\ell(q = 0) = 2\omega_0 \text{sign}(nH) + |\Omega_{ex}|(1 - \lambda_\ell/\lambda),$$

(22)

where $\ell = 1, 2 \ldots$ and $\lambda$ is related to the $\ell = 0$ case. The gapped modes are proportional to the polarization modified by the factor, depended on the interaction constants. Naturally, in the absence of the applied magnetic field, at $H = 0$, these gapped modes are related to the processes of the spin precession around the internal effective magnetic field originated from the electron exchange interactions.

For simplicity, we will calculate the wave-vector-dependent part of the first mode ($\ell = 1$) only. The corresponding dispersion equation of the gapped non-equilibrium spin waves is given as

$$\Omega_1(q = 0) = 2\omega_0 \text{sign}(nH) + |\Omega_{ex}|(1 - \lambda_1/\lambda),$$

(23)

and presented in fig. 3. The propagation of the gapped spin waves is also suppressed in the $|\Omega_{ex}| |\tau| < 1$ limit, while the relaxation time is reduced by the value of $\lambda_1/\lambda$ compared to the gapless spin wave in the paramagnetic regime, eq. (19). We also note, that the $q$-dependent part of the spectrum does not change sign at the ferromagnet-paramagnet transition contrary to the gapless mode.

**Amplification of the electromagnetic field.** – Consider the interaction of the electromagnetic field with the non-equilibrium spin-wave system. The energy of the field will be supplied to the system when the spin waves with positive frequency are excited. On the other hand, the energy is extracted from the system in the case of negative spin-wave frequency. In this regime the electromagnetic field interacting with the system might be amplified. Indeed, this statement can be verified by estimating the energy dissipation of the electromagnetic field in the spin system.

Experimentally, one might examine the absorption in the spin-wave resonance effect \[14,21\], in the vicinity of the spin-waves resonant frequency $\Omega$, where the wave vector $q$ is defined by the geometry of the sample. One obtains for the energy dissipation of the field

$$Q = \frac{1}{4\pi} \left\langle H(t) \frac{\partial B(t)}{\partial t} \right\rangle = |H|^2 \omega \Omega \left[ \chi^+(q, \omega) + \chi^-(q, \omega) \right].$$

(24)

The transverse susceptibility $\chi^+(q, \omega)$ is given by eq. (4) and can be written as

$$\chi^-(q, \omega) = -2\mu_\ell^2 \frac{N\mathcal{P}}{\omega - \Omega + i\gamma}.$$  

(25)

where $N$ is the electron concentration, $\mathcal{P} > 0$ is the spin polarization, while $\Omega$ is the real part of the spin-wave frequency given by eq. (13), $\gamma > 0$ is the damping of spin waves in the paramagnetic metal, see eq. (19), and in the ferromagnetic metal, eq. (14). The susceptibility $\chi^{-}(q, \omega)$ differs from eq. (25) by the opposite sign at $\omega + i\gamma$. The imaginary part in eq. (24) is

$$\Im \left[ \chi^+(q, \omega) + \chi^-(q, \omega) \right] = 2\mu_\ell^2 \frac{N\mathcal{P} \gamma}{(\omega - \Omega)^2 + \gamma^2} - \frac{N\mathcal{P} \gamma}{(\omega + \Omega)^2 + \gamma^2}.$$  

(26)

Finally, in the vicinity of the resonant frequency $\Omega$ we estimate the resulting expression for the energy dissipation

$$Q \sim \mu_\ell^2 |H|^2 \frac{N\mathcal{P} \Omega}{h\gamma}.$$  

(27)

The real part of the frequency $\Omega$ in the ferromagnet is always positive, eq. (14). Thus, the electromagnetic field contributes the energy to the spin system, which reveals in the well known spin-wave resonance effect, where one observes the microwave absorption peaks by the standing spin waves in thin ferromagnetic films. Note that the spin-waves' frequency is also positive in the vicinity of the paramagnet-ferromagnet phase transition $|1 - \lambda_\ell| \ll 1$ of the non-equilibrium paramagnet, fig. 2. In this regime the value of the dissipation, $Q > 0$, and the energy of the electromagnetic field is still supplied to the spin-wave system.

However, the energy influx to the magnetic system becomes negative, $Q < 0$, when the frequency changes sign, $\Omega \ll 0$, eq. (19), at $\lambda \nu < 1$. Therefore, the electron exchange interactions in the paramagnetic material with the non-equilibrium spin polarization can lead to the amplification of electromagnetic field. Experimentally, one might suggest the combination of the spin-wave resonance effect with the spin injection method. We propose the observation of the energy dips in the absorption spectrum instead of the expected absorption peaks in the spin-resonance effect.

**Conclusions.** – To conclude, we have studied the effect of electron exchange interactions on the non-equilibrium spin excitations in paramagnetic metals. We
have obtained the energy spectrum of the non-equilibrium spin waves in the diffusive paramagnetic system. Generally speaking, the negative imaginary part of the spin-wave frequency shows the stability of the non-equilibrium state under small spin-wave perturbations. In the regime when the electron-electron exchange energy is larger than the inverse electron mean free time the resulting spin waves are long-lived excitations, while in the diffusive limit the damping of spin waves is very strong. The excitation of these spin waves lowers the energy of the system due to the negative sign of the real part of the frequency. Interestingly, we have shown that the electromagnetic field can be amplified in the non-equilibrium spin-polarized paramagnetic metal. The spin injection plays the role of pumping source establishing the long-lived collective spin excitations which form an active medium leading to the amplification of the electromagnetic field.

***

We are grateful for the financial support of RFFI under Grant No. 10-02-00681-A and Federal Program under Grant No. 2009-1.5-508-008-012.

REFERENCES

[1] Dyakonov M. I. (Editor), Spin Physics in Semiconductors (Springer) 2008.
[2] Zutic I., Fabian J. and Das Sarma S., Rev. Mod. Phys., 76 (2004) 323.
[3] Meier F. and Zakharchenya B. P. (Editors), Optical Orientation (North-Holland, Amsterdam) 1984.
[4] Aronov A. G. and Pikus G. E., Sov. Phys. Semicond., 10 (1976) 698.
[5] Aronov A. G., JETP Lett., 24 (1976) 32.
[6] Johnson M. and Silsbee R. H., Phys. Rev. Lett., 55 (1985) 1790.
[7] Johnson M. and Silsbee R. H., Phys. Rev. B, 35 (1987) 4959.
[8] Aronov A. G., JETP, 73 (1977) 577.
[9] Berger L., Phys. Rev. B, 54 (1996) 9353.
[10] Słonczewski J. C., J. Magn. & Magn. Mater., 159 (1996) 1.
[11] Tsoi M., Jansen A. G. M., Bass J., Chiang W.-C., Seck M., Tsoi V. and Wyder P., Phys. Rev. Lett., 80 (1998) 4281.
[12] Sankey J. C., Cui Y.-T., Sun J. Z., Słonczewski J. C., Buhrman R. A. and Ralph D. C., Nat. Phys., 4 (2008) 67.
[13] Kubota H., Fukushima A., Yakushi K., Nagahama Taro, Yuasa S., Ando K., Maehara H., Nagamine Y., Tsunekawa K., Djayaprawira D. D., Watanabe N. and Suzuki Y., Nat. Phys., 4 (2008) 37.
[14] White R., Quantum Theory of Magnetism (Springer-Verlag) 1983.
[15] Moriya T., Spin Fluctuations in Itinerant Electron Magnetism (Springer-Verlag) 1985.
[16] Altshuler B. L. and Aronov A. G., Electron-Electron Interactions in the Disordered Conductors, edited by Efros A. L. and Pollak M. (Elsevier) 1985.
[17] Halperin B. I. and Hohenberg P. C., Phys. Rev., 188 (1969) 898.
[18] Fulde P. and Luther A., Phys. Rev., 170 (1968) 570.
[19] Platzman P. M. and Wolf P. A., Phys. Rev. Lett., 18 (1967) 280.
[20] van Dorpe P., Liu Z., van Roy W., Motyni V. F., Sawicki M., Borghs G. and De Boeck J., Appl. Phys. Lett., 84 (2004) 3495.
[21] Seavey M. H.-Jr. and Tannenwald P. E., Phys. Rev. Lett., 1 (1958) 168.