MAGNETIC MEDIA AS A SOURCE AND DETECTOR OF AXIONS

P.V. Vorobyev, I.V. Kolokolov and V.F. Fogel

Abstract

The new laboratory experiment for the generation and detection of axions using the coherent axion–magnon conversion in ferromagnets is discussed. The efficiency of the conversion process is ensured by the coherent axion-spin wave coupling(*).

1. Introduction

There are no doubts that the eventual discovery of a pseudo-scalar long-range force would drastically change our picture of the universe. The (pseudo) Goldstone bosons mediating such forces are massless (or very light) and can be detected by methods which can be considered as non-traditional in particle physics. Indeed, this problem may be formulated as that of detecting weak classical fields (static or oscillating), and thus the coherent interaction of such fields with macroscopic bodies can be used [1, 2]. In the experiment [1] the influence of a spin-polarized sample on the polarization of the others due to an intermediate (quasi-static) arion field (= massless axion field) has been measured. The result gave the limit $G_a < 10^{-3}\ G_F$, where G_F is the weak-interaction Fermi constant and G_a is the fermion–fermion interaction constant induced by the arion exchange.

There are proposals [3–6] for experiments on arion research, based on the coherent arion–photon conversion effect in a transverse magnetic field. Recently, the detailed project of the Sun axions detector has been published [7]. The coherence of axion–photon transformation in the case of massive axions is reached in ref. [7] by filling the interaction space with a light gas (H_2, He) ensuring the photon mass. Particles with mass up to 1 eV could be detected in such a way.

The fundamental Lagrangian of the axion (a) and fermion (f) interaction has the form

$$\mathcal{L}_{af} = q_f\,(\bar{f} i \gamma_5 f)\cdot a\ , \tag{1}$$

where q_f is the dimensionless axionic charge of the fermion. This vertex produces the effective axion–photon coupling

$$\mathcal{L}_{a\gamma} = \frac{1}{M}\, a\,(\vec{E}\cdot\vec{B}_0) \tag{2}$$

due to the well-known triangle anomaly of the axial current.

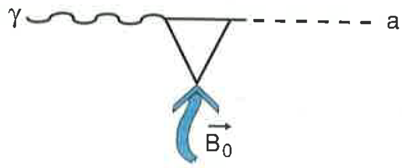

Diagram 1

Here $\vec{B}_0$ is the external magnetic field (denoted in eq. (2) by the double line), $\vec{E}$ is the photon electric field and M is the coupling constant of a mass dimension. Since both the photon and arion dispersion laws coincide, this linear coupling (eq. (2)) leads to the coherent arion–photon conversion (or, as is the case of ref. [7], axion–massive photon coherent conversion) in the presence of transverse $\vec{B}_0$.

However, in some theories the amplitude in diagram 1 vanishes. On the other hand, this graph is higher in order than the diagram corresponding to the following Compton process:

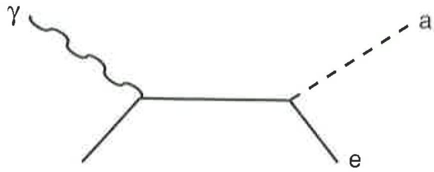

Diagram 2

It can be verified directly that this amplitude of arion–photon forward scattering vanishes when the electron is free. When the electron is in an external field, the first non-vanishing Compton diagram contains one quantum of that field.

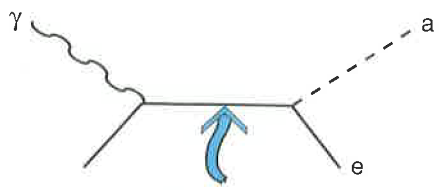

Diagram 3

(*) A short version of this work was published in Sov. Phys. JETP Lett. 50 (1989) 58.

Particle World, Vol. 1, No. 6, p. 163–167, 1990.

Here the double line denotes the external field seen as perturbation. This graph has the same number of electromagnetic vertices as the triangle of diagram 1. However, the three Compton amplitudes differ drastically from the ones associated to an anomaly graph such as diagram 1. When the electron is bound by an external field then the amplitude

$$\sim \sum_{n} \frac{|V_n|^2}{E - E_n} \qquad (3)$$

differs from zero owing to the presence of the field. Here, the solid line denotes the propagator of the bound electron, and the summation runs over the intermediate states with energies E_n and V_n, which are the corresponding matrix elements. The electron binding energy has an electromagnetic origin, and thus the energy denominator in the latter amplitude compensates the factor α in the matrix elements V_n when the energy of the initial particle is less than this binding energy.

In this range of energies the axions interact effectively only with the electronic magnetic moment density $\vec{M}(\vec{r})$. The Lagrangian eq. (1) is reduced to

$$\mathcal{L}_{ma} = x\,\vec{\nabla} a\ \vec{M}(\vec{r})\ , \qquad (4)$$

where $x = \mu_a/\mu_B = q_e/2e$, $\mu_a = q_e/2m_e$ being the arionic magneton and μ_B the Bohr magneton.

Here we consider the case where the arions only excite the spin degrees of freedom. For our purpose it is convenient to use the classical description of the arion field and spin waves. The notes cited above may be regarded as the interpretation of the final results, in terms of the fundamental processes.

2. The ferromagnetic detector of (pseudo) Goldstone bosons

Let us consider the detector schematically pictured in fig. 1. Contrary to the laser one discussed in refs [3] and [4], here the space between the magnetic poles is filled with a non-conducting ferromagnet. In the detector the microwave rather than optical diapason is used. The frequencies of the pumping generate a coherent arion field. The narrow arion beam passes freely through the walls of the waveguide as well as the screens that absorb photons and comes into the second waveguide equally filled with magnetized ferrite. The arion wave excites resonantly the

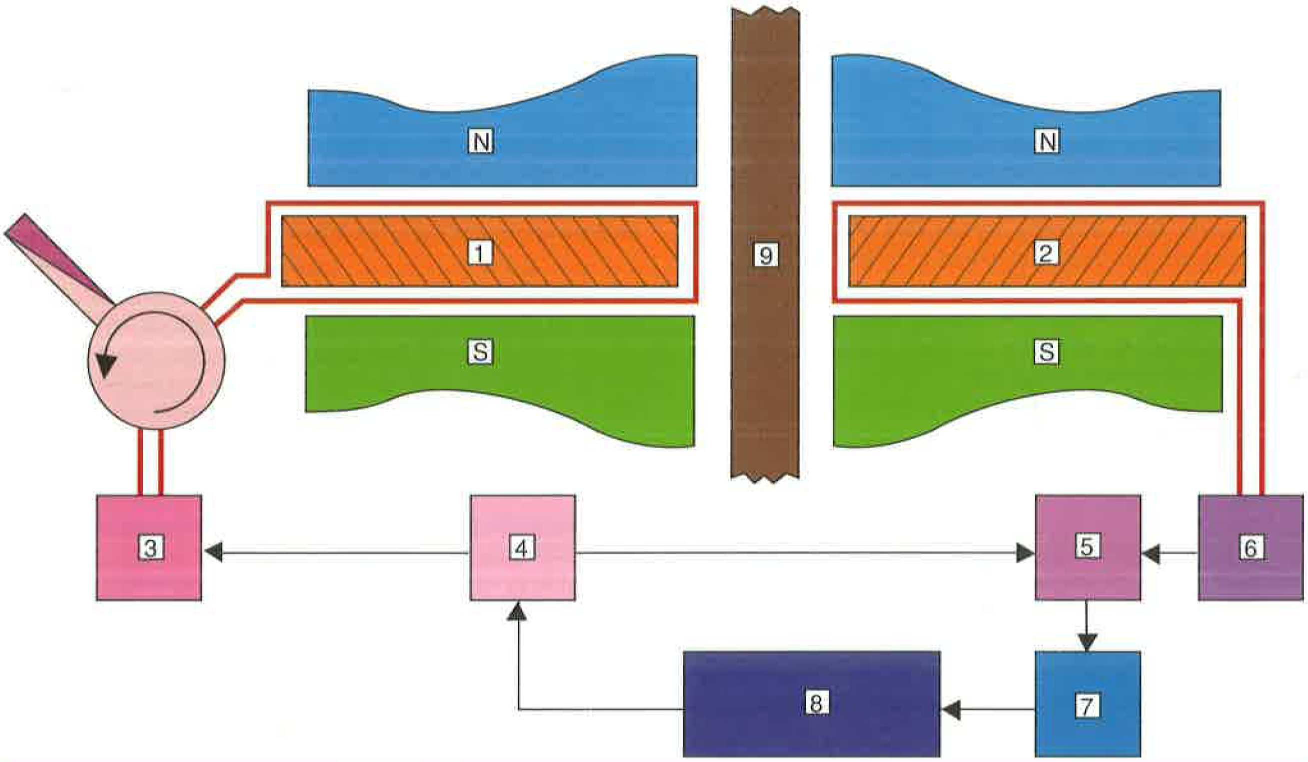

FIGURE 1

Scheme of the axion detector: (1) and (2) waveguides filled with ferrite; (3) powerful high-frequency generator; (4) synchrogenerator; (5) synchrodetecto; (6) sensitive receiver; (7) ADC; (8) computer; and (9) screen.

coherent spin precession and the weak electromagnetic wave coupled to the spins. The latter wave is registered by a high-sensitivity receiver (e.g. maser or Rydberg atom receiver).

The second (detecting) resonator may be either open (i.e. a long waveguide loaded on the receiver) or closed. In the former case the spin-wave amplitude on the exit of the waveguide is enhanced by the factor kL, where k is the wave vector and L is the waveguide length. The response amplitude of the closed resonator is proportional to the quality factor ω/γ, where ω and γ are the frequency and the relaxation constant of the corresponding mode. The excitation of the eigenmode in the closed resonator may be considered as the result of iterated reflections on the relaxation time interval. In other words, the closed resonator is equivalent to an open one having an effective length $L_{eff} = \omega/k\gamma$. In our frequency range the inequality $\omega/\gamma >> kL$ holds for lengths of the order of one metre and it is thus more convenient to use the closed resonator (this is not the case when the sharp radiation pattern is needed for the receiver).

Let us consider the operating principle and the possibilities of such a detector in more detail.

First, we consider the interaction of a pseudoscalar arion field with the electromagnetic one in the presence of a paramagnetic or ferromagnetic medium magnetized up to saturation.

The dynamical variable of the medium is the field $\vec{m}(\vec{r})$ of the magnetization density variations. The Lagrangian of the magnet interaction with the time-dependent magnetic field $\vec{b}(\vec{r})$ is equal to

$$\mathcal{L}_{\rm mem} = \vec{b}(\vec{r})\,\vec{m}(\vec{r}) \ . \qquad (5)$$

It follows from eq. (4) that the coupling of the magnet to the arion field can be described by the Lagrangian

$$\mathcal{L}_{\rm ma} = x\,\vec{\nabla} a \bullet \vec{m} \ , \quad x = \mu_{\rm a}/\mu_{\rm B} \ . \qquad (6)$$

The proper dynamics of the field $\vec{m}(\vec{r})$ are governed by the Bloch equation

$$\dot{\vec{m}}(\vec{r}) = 2\mu_{\rm B}/\hbar\left[\delta H/\delta\vec{M}(\vec{r}), \vec{M}(\vec{r})\right] , \qquad (7)$$

where H is the magnet's Hamiltonian, $\vec{M}(\vec{r}) = \vec{M}_0 + \vec{m}(\vec{r})$, $\vec{M}_0$ being the equilibrium magnetization directed along the z-axis.

It is more convenient to analyse the magnet dynamical equations in terms of the canonical variables C, C^* [8]

$$i\dot{C} = \delta H/\delta C^* \ , \qquad (8)$$

connected with the magnetization density by the classical version of the Holstein–Primakoff transformation

$$m_+ = m_{\rm x} + i m_{\rm y} = C\sqrt{\omega_{\rm m}\left(1 - \frac{C^*C}{2M_0}\right)}, \quad m_{\rm z} = -gC^*C \ . \qquad (9)$$

Here, $\omega_{\rm m} = 2gM_0$, $g = 2\mu_{\rm B}/\hbar \approx 2\pi \bullet 2.8$ MHz/$G_{\rm s}$. Neglecting the non-linearity for the case $m_{\rm x}$, $m_{\rm y} << M_0$ we have

$$m_+ = C\sqrt{\omega_{\rm m}} \ , \quad m_{\rm z} = -gC^*C \ . \qquad (10)$$

The exchange interaction is unessential at the frequencies involved (it can, however, be taken into account easily) and the eigenmodes of the resonator filled with magnet can be determined by solving the Maxwell–Bloch equations derived from the Lagrangian

$$\mathcal{L} = -\frac{1}{16\pi}F_{\mu\nu}^2 + \vec{m}\vec{b} + i\dot{C}C^* - \Omega_0 C^*C \ . \qquad (11)$$

Here, $\Omega_0 = gB_0$, $\vec{B}_0$ is the constant homogeneous part of the true magnetic field in the medium, and $\vec{b}$ is the dynamical variable of the electromagnetic field: $F_{\rm ij} = \varepsilon_{\rm ijk} b_{\rm k}$. It is worth noting that the magnetic dipole–dipole interaction is automatically taken into account by the Lagrangian of eq. (11). Restricting ourselves to the case of the infinite in the (y–z) medium plane we consider the fields depending on x alone. The Lagrangian of the pumping field-spin wave coupling in terms of variables C, C^* has the form

$$\mathcal{L}_{\rm mem} = \vec{m}\vec{b} = \frac{ib}{2}\sqrt{\omega_{\rm m}}\left(C - C^*\right) . \qquad (12)$$

The linear polarization of the pumping wave along the axis is assumed [i.e. $\vec{b} = (0, b, 0)$]. The eigenvectors of the coupled oscillations of the magnetic moment and the electromagnetic field can be found from the linearized Maxwell–Bloch equations. For the waves propagating only along the x-axis, we have

$$(\varepsilon\partial_{\rm t}^2 - \partial_{\rm x}^2)b = -2\pi i\,\partial_{\rm x}^2\sqrt{\omega_{\rm m}}\left(C - C^*\right) ,$$

$$i\dot{C} - \Omega_0 C = -i/2\sqrt{\omega_{\rm m}}\,b \qquad (13)$$

(here ε is the dielectric constant of the medium). From this system of equations we obtain the structure of eigenmodes:

$$C_{\rm k} = \left(A\,e^{-i\Omega t} + D\,e^{i\Omega t}\right)\sin kx \ ,$$

$$B_{\rm k} = \left(B\,e^{-i\Omega t} + B^* e^{i\Omega t}\right)\sin kx \ ,$$

$$A = -\frac{i}{2}\frac{\sqrt{\omega_{\rm m}}}{\Omega - \Omega_0}B \ , \quad D^* = -\frac{i}{2}\frac{\sqrt{\omega_{\rm m}}}{\Omega + \Omega_0}B \qquad (14)$$

and the dispersion law

$$\Omega_{\rm k}^2 = \frac{1}{2\varepsilon}\left[\left(k^2 + \varepsilon\Omega_0^2\right) \pm \sqrt{\left(k^2 - \varepsilon\Omega_0^2\right)^2 + 4\varepsilon k^2 \Omega_{\rm M}\Omega_0}\,\right] . \qquad (15)$$

This is illustrated by the two curves in fig. 2. The oscillations proper to the magnetic moment $\vec{m}(\vec{r})$ are described by the modes corresponding both to the horizontal asymptotics of the lower branch of the dispersion law and to the $\Omega \sim \Omega_0$ domain of the upper branch. Since the inequality $\varepsilon > 1$ holds, the linear arion dispersion law $\Omega_k = k$ always has an intersection point with the upper curve. Thus, the arion field can be excited in a space-coherent way.

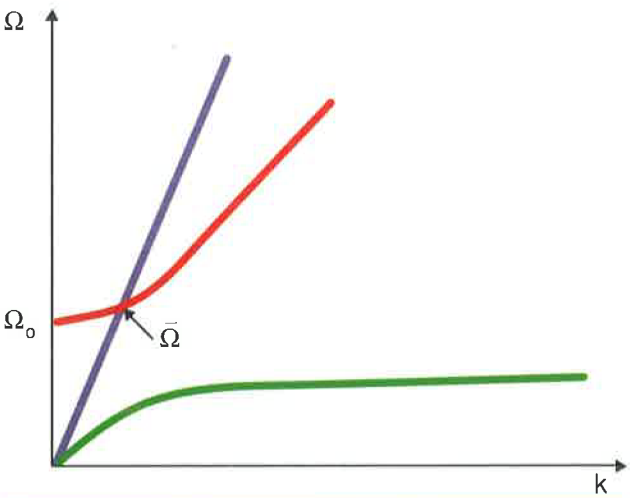

FIGURE 2

Dispersion laws of the magnon–photon coupled oscillations; the asymptotics of the lower branch is $\sqrt{\Omega_0(\Omega_0 - \Omega_M)}$, for the upper branch it is $\Omega_k = k/\sqrt{\varepsilon}$. The arion dispersion line $\Omega_k = k$ intersects the upper curve at the point $\Omega = \bar{\Omega}$.

Since the arion field interacts only with the spin waves, the corresponding amplitude in the excited mode ought to be quite large. This means that the intersection point of the line $\Omega_k = k$ and the upper dispersion curve should have a frequency $\bar{\Omega}$ comparable with the Ω_0. On the other hand, this frequency $\bar{\Omega}$ should not be too close to Ω_0 because the group velocity of the wave tends to become zero in the limit $\bar{\Omega} \to \Omega_0$ and the signal could be lost. The proper tuning can be reached by a suitable choice of the external magnetic field and it is easily realized at frequencies $\bar{\Omega} \cong 2\pi \cdot 10^{10}$ Hz.

The amplitude of the mode (eq. (14)) in the generating resonator is defined by the energy balance (see also below): if $\mathcal{E}$ is the energy per unit of time coming into this resonator, then, neglecting the pure electromagnetic energy loss, we have

$$\frac{1}{2} \cdot \left(|A|^2 + |D|^2\right) = \mathcal{E} / (\gamma \Omega_0 \cdot \text{volume}) \qquad (16)$$

(here γ is the magnon relaxation frequency). The arion field generation is governed by eq. (6) having in the variables C, C^* the form

$$\Box a = -\frac{x}{2} \sqrt{\omega_m}\, \partial_x \left(C + C^*\right) . \qquad (17)$$

If the excited mode (eq. (14)) has the frequency $\bar{\Omega}$, then the arion field amplitude at the output of the first resonator is equal to

$$a = i \frac{x}{2} L \sqrt{\omega_m} \left[A \exp\left(i\bar{\Omega} t\right) + D \exp\left(-i\bar{\Omega} t\right)\right] . \qquad (18)$$

Here L is the length of the resonator.

The same mode (eq. (14)) is excited resonantly by the arion wave (eq. (18)) in the second resonator separated from the first one by the screen. The resulting amplitude D_f is expressed in terms of the amplitude D (eq. (14)) as follows:

$$D_f = D \cdot \frac{\omega_m x^2 (kL)}{4\gamma} ,$$

in accordance with the equation of motion

$$i\dot{C} - \Omega_0 C + \frac{i}{2} \sqrt{\omega_m}\, b = \frac{x}{2} \sqrt{\omega_m}\, \partial_x a . \qquad (19)$$

The square of the amplitude of the electromagnetic oscillations coupled with the magnetic moment is equal to

$$|B_f|^2 = x^2 \omega_m (kL)^2 |B|^2 \frac{\left(\Omega_0 - \bar{\Omega}\right)^2}{16\gamma^2} . \qquad (20)$$

Taking into account the equality (16), we come to the expression for the relation between the energy flow at the end of the second resonator $\mathcal{E}_f$ and the pumping energy flow $\mathcal{E}$

$$\mathcal{E}_f / \mathcal{E} \approx (kL) x^4 \frac{\bar{\Omega} \omega_m \left(\Omega_0 - \bar{\Omega}\right)^2}{8\gamma^3 \bar{\Omega}} \qquad (21)$$

(the amplitude A is neglected in comparison with D).

The value of expression (21) can be estimated taking all the frequencies to be of the same order

$$\mathcal{E}_f / \mathcal{E} \cong x^4 (kL) \left(\bar{\Omega}/\gamma\right)^3 . \qquad (22)$$

For normally used ferromagnets, the factor $\bar{\Omega}/\gamma$ has an order of magnitude of $\sim 10^3$ and for $kL = 10^2$ we have

$$\mathcal{E}_f / \mathcal{E} \approx 10^{11} x^4 . \qquad (23)$$

The direct double photon–arion transformation has the conversion coefficient

$$\mathcal{E}_f / \mathcal{E} \approx \alpha^4 x^4 (kL)^4 , \qquad (24)$$

which is about 11 orders of magnitude less than expression (23) with the same parameter values.

It is worth noting some important issues. The spin-wave amplitude in the generating resonator is limited by the thresholds of instabilities (see below). Increasing the quality of this resonator is of no advantage whilst the increase in length is more profitable. Thus, it is reasonable to use, in the first resonator, a sample magnet of sufficient length. As a rule, the quality factors of such ferrites do not exceed 100. On the other hand, the amplitude of the spin wave excited by arions (axions) in the detecting resonator is proportional to $(\bar{\Omega}/\gamma)^2$ and, additionally, the question of the threshold does not arise here. Consequently, the magnetic sample in the detecting resonator should have the quality $\bar{\Omega}/\gamma$ as high as possible. Monocrystals from the yttrium–iron garnet (YIG) having the well-known spin-wave dynamics [9, 10] have qualities of the order of 10^4, at the frequencies near the bottom of the spin-wave spectrum, and it is worthwhile using such a magnet as the detector.

3. Non-linear limits

With the increase of the pumping power, the spin-wave amplitude in the generating resonator reaches the threshold value; then the spin wave becomes unstable and any further increase of the pumping power has no effect. The review quoted in ref. [8] is devoted to the non-linear theory of spin waves and the estimations given below are taken from that book.

The lowest threshold has the decay instability: the spin wave transforms into two waves with half frequencies. The threshold amplitude has the value

$$C_{\rm th} \approx \frac{\gamma}{g\bar{\Omega}} \sqrt{\omega_{\rm m}} \qquad (25)$$

and corresponds to the absorbing power $\mathcal{E} \approx V \cdot \gamma^3/2\pi g^2$. For a YIG sample with volume $V = 100$ cm^3 we have 10 W. However, this instability develops when the decay conservation laws are fulfilled. If $\bar{\Omega} \approx \Omega_0$, our excitation can transform into two magnons of the lower branch only. When the frequency $\Omega_0/2$ belongs to the forbidden zone

$$\Omega_0 > 2\sqrt{\Omega_0(\Omega_0 - \Omega_{\rm M})}\ , \qquad (26)$$

then this transformation is suppressed (more precisely: in this case the decay products belong to the exchange part of the spectrum, where the relaxation constants γ_k are sufficiently greater than in the magnetostatic one; it is worth noting that the γ in formula (25) is only the secondary waves relaxation frequency). This inequality solved with respect to Ω_0 has the form

$$\Omega_0 < 4/3\,\Omega_{\rm M}\ . \qquad (27)$$

On the other hand, the stability of the sample, with respect to the formation of the domain structure, demands $\Omega_0 > \Omega_{\rm M}$. Thus, to avoid these instabilities we should be restricted by the frequencies diapason

$$\Omega_{\rm M} < \Omega_0 < 4/3\,\Omega_{\rm M}\ . \qquad (28)$$

The spin-wave scattering processes are always permitted and the corresponding modulation instability has the threshold

$$C_{\rm th}^{(2\to 2)} \cong \frac{1}{g}\sqrt{\frac{\gamma}{\Omega^{(*)}}}\sqrt{\omega_{\rm m}}\ . \qquad (29)$$

The absorbed energy power reached in this case is $\bar{\Omega}/\gamma$ times greater than in the former case of decay instability.

4. Conclusions

Our method is based on the coherent arion–magnon transformation provided by the intersection of the dispersion curves of that excitation. Thus, the mass of the detected axion does not exceed the gap in the magnon spectrum, i.e. $\sim 10^{-4}$ eV.

Our axion detector could be used for the search of cosmic axions. One of its advantages is the possibility of fine-tuning the frequency by changing the external field.

We are grateful to A.A. Anselm, I.B. Khriplovich and V.B. Cherepanov for stimulating discussions and helpful advice.

References

[1] P.V. Vorobyev and Ya.I. Gitarts, Phys. Lett. 208 (1988) 146.
[2] A.A. Anselm, Proc. of the XXth Winter School of Leningrad, Nucl. Phys. Inst., Leningrad, USSR (1988).
[3] A.A. Anselm, Phys. Rev. D37 (1988) 2001.
[4] K. Van Bibber et al., Phys. Rev. Lett. 59 (1987) 759.
[5] P. Sikivie, Phys. Rev. D32 (1983) 2988.
[6] A. Melissinos et al., Brookhaven proposal, AGS P840 (1988).
[7] K. Van Bibber et al., Phys. Rev. D39 (1989) 2089.
[8] V.S. L'vov, Non-linear theory of spin waves, Moscow, USSR (1987) (in Russian).
[9] I.V. Kolokolov, V.S. L'vov and V.B. Cherepanov, Sov. Phys. JETP 86 (1984) 1946.
[10] A.H. Anisimov and A.G. Gurevich, Sov. Phys. JETP 68 (1976) 677.

Address:
USSR Academy of Science
Siberian Division
Institute of Nuclear Physics
630090 Novosibirsk, USSR

Received on 10 June 1990.