Antiferromagnetic axion detector

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Antiferromagnets with easy-plane anisotropy are proposed as axion detectors. It is shown that the response of the detector is proportional to the ratio of the Dzyaloshinskiĭ field to the external magnetic field; this makes it possible to strengthen the bounds on the axion-electron interaction constant.

1. INTRODUCTION

The discovery of long-range action transferred by massless or almost massless pseudoscalar particles (arions, axions) would be of great significance both for constructing an adequate cosmological model and for studying physics at very small distances. The experimental status of the search for such "exotic" long-range actions at the beginning of 1989 is reviewed in Ref. 1.

The interaction of an axion field a with fermions ψ in matter is described by the Lagrangian

$$\mathscr{L}_{int} = -qa\bar{\psi}i\gamma^5\psi, \qquad (1.1)$$

where q is the axion charge of the fermion.

We shall study the interaction of the field a with magnetically ordered dielectrics. At sufficiently low energies an axion can excite only the spin degrees of freedom of the electrons localized at the sites of the crystal lattice. In this limit the effect Lagrangian, following from Eq. (1.1), coupling the axion field with the medium is equal to

$$\mathscr{L}_{ma} = \varkappa \nabla a \mathbf{m}(\mathbf{r}). \qquad (1.2)$$

Here $\varkappa = \mu_a/\mu_B$ is the ratio of the axionic magneton of the electron to the Bohr magneton, $\mu_a = q/2m_e$, m_e is the electron mass, and $\mathbf{m}(\mathbf{r})$ is the magnetization density of the medium. [The relation $m_e\bar{\psi}i\gamma^5\psi = \partial_\mu\bar{\psi}\gamma^5\gamma^\mu\psi +$ the contribution of the axial anomaly, which we ignored and which leads to direct axion-photon conversion in an external field[2,3] (see also Ref. 1), was used in the derivation of the Lagrangian (1.2).]

This "quasimagnetic" character of the field ∇a implies, in particular, that the static axion field generated by a ferromagnet induces a constant magnetization in paramagnetic samples separated from this ferromagnet by a superconducting screen.[4] The experiment of Ref. 4, performed according to this scheme, gave a limit on the constant $\varkappa$:

$$\varkappa < 2\cdot 10^{-7}. \qquad (1.3)$$

In Refs. 5 and 6, on the other hand, an experiment on the generation and detection of a dynamic axion field was studied. This experiment employed the existence of a point of intersection of the dispersion relations of axions and spin waves in a ferromagnet. Detailed calculations of the coefficients of coherent axion-magnon and double magnon-axion-magnon conversion in a ferromagnetic medium were presented in Ref. 6. (In Ref. 7, which appeared at the same time as Ref. 5, it is suggested that galactic axions with wavelength 10–100 m be detected based on their resonance interaction with the homogeneous (ferromagnetic precession.) Magnons excited in the ferromagnetic detector[5,6] are detected based on the electromagnetic oscillations which are coupled with them. Ultimately it is precisely the number of these coupled photons, which is proportional to $\varkappa^4$, that is recorded.

The number of excited magnons can be determined independently by measuring the variation of the macroscopic magnetic field (which is proportional to the measurement of the macroscopic magnetization) with the help of a SQUID magnetometer. It is known, however, that antiferromagnets (more accurately, weak ferromagnets) with anisotropy of the "easy plane" type and large Dzyaloshinskiĭ field H_D are more suitable for such measurements (see, for example, Ref. 8). In this case the decrease in the weak ferromagnetic moment accompanying excitation of one (quasi) goldstone magnon can exceed μ_B by more than an order of magnitude (this phenomenon was used in Ref. 8).

In this paper we shall study the possibility of such an antiferromagnetic detector of axions. We shall show that the response to the axion signal measured by a SQUID contains the enhancement factor H_D/H_0, where H_0 is the external magnetic field. An experiment of the type described in Ref. 5 on the generation and detection of axions can also strengthen the limit on the constant $\varkappa$ by two orders of magnitude.

2. AXION-MAGNON CONVERSION IN AN ANTIFERROMAGNETIC MEDIUM

Let $\mathbf{M}_1$ and $\mathbf{M}_2$ be the magnetizations of the sublattices of the antiferromagnet,

$$\mathbf{M}_1 = \mathbf{M}_0 + \mathbf{m}_1, \quad \mathbf{M}_2 = -\mathbf{M}_0 + \mathbf{m}_2,$$

where $2\mathbf{M}_0$ is the (almost) equilibrium value of the antiferromagnetic moment and $\mathbf{m}_{1,2}$ are the dynamic variables of the medium. Let the z-axis be oriented along $\mathbf{M}_0$ and $|\mathbf{m}_{1,2}| \ll |\mathbf{M}_0|$. Then the transformation to canonical variables c_1, c_2 (see, for example, Ref. 9), neglecting higher-order nonlinear terms, is made by means of the following substitution:

$$\begin{gathered} m_{1x} + im_{1y} = c_1(2\omega_m)^{1/2}\left(1 - \frac{g}{4M_0}c_1^*c_1\right), \\ m_{2x} - im_{2y} = c_2(2\omega_m)^{1/2}\left(1 - \frac{g}{4M_0}c_2^*c_2\right), \\ \omega_m = gM_0, \quad g = 2\mu_B/\hbar \approx 2\pi\cdot 2{,}8 \text{ MHz/Oe}, \\ m_{1z} = -gc_1^*c_1, \quad m_{2z} = gc_2^*c_2 \end{gathered} \qquad (2.1)$$

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(the classical analog of the Holstein-Primakov transformation). The isotropic exchange interaction density has the form

$$\mathscr{H}_{ex}=\alpha\mathbf{M}_1\mathbf{M}_2+\beta\partial_i\mathbf{m}_1\partial_i\mathbf{m}_2. \tag{2.2}$$

For wavelengths of interest to us the contribution of the nonuniform exchange can be neglected, so that $\mathscr{H}_{ex}\approx\alpha\mathbf{M}_1\mathbf{M}_2$.

In this paper we study crystals with local anisotropy of the easy plane type. For our choice of coordinate axes the energy density of such anisotropy is equal to

$$\mathscr{H}_a=\frac{\Omega_a}{\omega_m}(m_{1x}^2+m_{2x}^2), \tag{2.3}$$

where Ω_a is the corresponding frequency. Finally, the Dzyaloshinskiĭ interaction Hamiltonian density has the form

$$\mathscr{H}_D=\frac{H_D}{M_0}(M_{1z}M_{2y}-M_{2z}M_{1y}). \tag{2.4}$$

Making the substitution (2.1) in the total Hamiltonian

$$\mathscr{H}=\mathscr{H}_{ex}+\mathscr{H}_a+\mathscr{H}_D+\mathscr{H}_0,$$

we can see that $\mathscr{H}_D$ contains terms that are linear in $c_{1,2}$. Hence the configuration $c_1=c_2=0$ is not a vacuum configuration. The transfer to variables describing oscillations of the system near the true equilibrium is made, as is easily verified, by means of the translation

$$c_1\to c_1+i\beta,\quad c_2\to c_2-i\beta, \tag{2.5}$$

$$\beta=-\frac{1}{2}\left(\frac{\omega_m}{2}\right)^{1/2}\frac{H_D+H_0}{\alpha\omega_m}.$$

This corresponds to noncollinearity of the sublattices in the ground state and the appearance of a weak ferromagnetic moment along the y-axis. The terms in the Hamiltonian which are cubic in the starting variables will make, after the translation (2.5), a contribution to the quadratic part, which we write out in the final form as follows:

$$\begin{aligned}\mathscr{H}=&\alpha\omega_m(c_1^*c_1+c_2^*c_2+c_1c_2+c_1^*c_2^*)\\&+\Omega_a(c_1^*c_1+c_2^*c_2+{}^1/_2(c_1^2+c_2^2+\text{c.c.}))\\&+\alpha\beta^2g^2\{4(c_1^*c_1+c_2^*c_2)+2(c_1c_2+\text{c.c.})-{}^1/_4(c_1^2+c_2^2+\text{c.c.})\\&-{}^5/_2(c_2^*c_1+\text{c.c.})\}-4\alpha g^2\beta\tilde{\beta}((c_1^*c_1+c_2^*c_2)\\&+(c_1c_2+\text{c.c.})-(c_1^*c_2+\text{c.c.})).\end{aligned} \tag{2.6}$$

Here c.c. denotes the complex-conjugate terms, and

$$\tilde{\beta}=-\frac{1}{2}\left(\frac{\omega_m}{2}\right)^{1/2}\frac{H_0}{\alpha\omega_m}. \tag{2.7}$$

The Hamiltonian density (2.6) is diagonalized by canonical transformations in two stages (see, for example, Ref. 9): First the substitution

$$c_1=2^{-1/2}(d_1+d_2),\quad c_2=2^{-1/2}(d_1-d_2) \tag{2.8}$$

puts $\mathscr{H}(c_1,c_2)$ into the form $\mathscr{H}_1(d_1)+\mathscr{H}_2(d_2)$. Then $\mathscr{H}_1$ and $\mathscr{H}_2$ are each diagonalized by their own Bogolyubov transformation. The amplitude d_1 corresponds to the lowest mode of the oscillations which is excited under our conditions. For it we have

$$d_1=ub+vb^*,\quad \mathscr{H}=\omega_0b^*b, \tag{2.9}$$

$$\omega_0=[\Omega_0(\Omega_0+\Omega_D)]^{1/2},$$

$$\Omega_0=gH_0,\quad \Omega_D=gH_D,$$

$$u\approx\frac{1}{2\xi^{1/4}}(1+\xi^{1/2}),\qquad v=-\frac{1}{2\xi^{1/4}}(1-\xi^{1/2}),$$

where

$$\xi\approx\frac{2g^2\beta\tilde{\beta}}{\omega_m}=\frac{H_0(H_0+H_D)}{4\alpha^2M_0^2}\ll1. \tag{2.10}$$

We use the ratio of the parameters $\alpha\omega_m\gg\Omega_0$ and $\alpha\omega_m\gg\Omega_D\gg\Omega_0$, which are realized in external fields that are not too strong. (For example, in $FeBO_3$ we have $H_D\approx100$ kG.)

We present an expression, which follows from Eq. (1.2), for the Hamiltonian characterizing the interaction of the field b and the axion field for the case of an axion wave propagating along the z-axis:

$$V=i2^{1/2}g\beta\varkappa\partial_za(u-v)(b-b^*)\approx i2^{1/2}g\beta\varkappa\xi^{-1/4}\partial_za(b-b^*). \tag{2.11}$$

It can be shown that for all other directions of propagation of axions the response of the system will be suppressed.

From Eq. (2.11) it follows that a monochromatic axion wave with frequency ω_0 and amplitude a_0 excites linearly a spin wave with amplitude

$$b_0=i2^{1/2}g\beta\varkappa\frac{\omega_0}{\gamma}\xi^{-1/4}a_0, \tag{2.12}$$

where γ is the relaxation rate of the spin wave. The change in the static ferromagnetic moment has the following form in the variables c_1 and c_2:

$$\Delta M_y=-\frac{\beta}{4M_0}g\left(\frac{\omega_m}{2}\right)^{1/2}[4(c_1^*c_1+c_2^*c_2)-(c_1^2+c_2^2+\text{c.c.})]. \tag{2.13}$$

Assuming that only the lower mode with the amplitude (2.12) is excited and keeping only the resonance terms we obtain

$$\begin{aligned}\Delta M_y&\approx-\frac{\beta}{M_0}g\left(\frac{\omega_m}{2}\right)^{1/2}(u^2+v^2-uv)b_0^*b_0\\&=\frac{3}{2}\left(\frac{\omega_m}{2}\right)^{1/2}\frac{\beta^3}{\xi M_0}g^3\varkappa^2a_0^2\left(\frac{\omega_0}{\gamma}\right)^2\\&\approx\frac{3}{16}\varkappa^2(ga_0)^2\left(\frac{\omega_0}{\gamma}\right)^2\frac{H_D}{H_0}\frac{\Omega_D}{\alpha\omega_m}M_0.\end{aligned} \tag{2.14}$$

By reducing the external field it is possible to make the ratio H_D/H_0 assume a significant value and, therefore, to increase the variation of the static field $4\pi\Delta M_0$.

The linear response is enhanced for the z-component of the magnetization. The amplitude m_z^0 is proportional to $\xi^{-1/2}$:

$$m_z^0\propto g^2\beta^2\varkappa\frac{\omega_0}{\gamma}\xi^{-1/2}a_0. \tag{2.15}$$

In unbounded space, owing to its longitudinal character, such a wave of oscillations of m_z will not generate an electromagnetic field. In a closed resonator with dimensions of the order of the wavelength, however, the overlap integral of the electromagnetic mode and the longitudinal spin wave (2.15) can be of order unity. Correspondingly the energy density of the coupled photons per unit volume is

$$\varepsilon_\gamma\sim\varkappa^2(a_0g)^2\left(\frac{\omega_0}{\gamma}\right)^2\frac{H_D}{H_0}\left(\frac{\Omega_D}{\alpha\omega_m}\right)^2M_0^2. \tag{2.16}$$

3. CONCLUSIONS

An antiferromagnetic detector can be employed to search for relic axions having a temperature of several Kelvins. Another application is in experiments of the type proposed in Ref. 5: A spin wave with a large amplitude in one resonator generates an axion wave which passes freely through a superconducting screen and excites another spin wave in a second (detecting) resonator. If, as one can see from the relation (2.15), an antiferromagnet can be more advantageous as a detector, then it is preferable to use a ferromagnet in the generator. Indeed, spin waves in antiferromagnets are much more highly nonlinear[9] and different instabilities develop at significantly smaller amplitudes. SQUID measurements of ΔM_y with a sensitivity of 10^{-11} $G\cdot cm^2$ are unlikely to strengthen appreciably the limit (1.3) on $\varkappa$ and can be regarded only as an independent experiment with the given accuracy. It is of greater interest to use closed resonators and Rydberg photon counters.

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