Coherent axion-photon transformations in the forward scatterings on atoms

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In certain laboratory experiments the production and/or detection of axions is due to the photon-axion transformations in a strong magnetic field. This process is coherent, and the rate of the transformation is proportional to the length $l$ and magnitude $B$ of the magnetic field squared, $\sim l^2 B^2$. In the present paper, we consider coherent production of axions due to the forward scattering of photons on atoms or molecules. This process may be represented as being due to an effective electromagnetic field which converts photons to axions. We present an analytical expressions for such effective magnetic and electric fields induced by resonant atomic M0 and M1 transitions, as well as give some numerical estimates for these fields. The corresponding experiments would allow to measure the electron-axion coupling constant $g_{ae}$ in the same way as the photon-axion coupling $g_{a\gamma}$ is studied.

I. INTRODUCTION

The axion was introduced to explain the absence of CP violation in the strong interaction [1–9] and is considered as a probable candidate for a dark matter particle. The majority of experiments searching for axions in laboratories are based on the transformation of axions to photons (for axion detection) and photons to axions (for axion production) in a strong magnetic field [10–14]. The study of atomic and molecular transitions for the axion detector has also been suggested in Refs. [50–55]. It is natural to assume that the atomic transitions can play role of effective magnetic and electric fields stipulated by the interaction of photons on atoms near atomic resonances.

The photon-to-axion conversion probability in magnetic field $B$ is [56] (in units $\hbar = c = 1$)

$$P = \frac{\omega}{4k_a} (g_{a\gamma}B)^2 F^2(q),$$  

where $\omega$ is the axion and photon energy, $k_a$ is the axion momentum, $g_{a\gamma}$ is the axion-photon coupling constant, $l$ is the spatial extent of the magnetic field in the direction of the axion-photon beam, $F(q) = \int dx \, e^{-iqx} B(x)/B$ is the form factor, $q = \omega - k_a$ is the momentum transferred from the photon field to the axion field. For the axion mass $m_a \ll \omega$, we have $q \approx 0$ and $F(0) \approx 1$.

Probabilities of the transformation of photons to axions and axions to photons are calculated by solving the system of Maxwell and Klein-Gordon equations coupled through the interaction term in the Lagrangian density:

$$L^{a\gamma} = g_{a\gamma} \, a \, (E \cdot B),$$  

where $a$, $E$ and $B$ are the axion, electric and magnetic fields, correspondingly. A similar effective Lagrangian density appears due to the interaction between the axion field and an atom

$$L^{\gamma a}_{\text{eff}} = g_{ae} \, a \, (E \cdot B_{\text{eff}} + E_{\text{eff}} \cdot B),$$

where the effective magnetic $B_{\text{eff}}$ and electric $E_{\text{eff}}$ fields are produced by the atom, while $E$ and $B$ are physical electromagnetic fields of an external photon. The effective Lagrangian (3) includes the coupling constant $g_{ae}$ which originates from the axion-electron (or axion-nucleon) interaction

$$L^{ae} = g_{ae} \, \bar{\psi} \gamma^\mu \gamma^\nu \gamma^\rho \psi,$$

where $\psi$ is the electron (or nucleon) spinor field.

In this note, we derive the expressions for the effective electric and magnetic fields stipulated by the interaction (4).

II. CALCULATIONS

Allowing for the axion-electron interaction [1], one can consider a scattering amplitude, where a bound electron absorbs an incident photon and emits an axion. Of special interest is the scattering in forward direction since the forward scattering on multiple atoms is always coherent. If the wavelengths of the photon and axion are significantly larger than the distance between the atoms, we may treat the atomic system as a continuous medium, and the forward scattering amplitude scales linearly with the number density of atoms $N_{\text{at}}$. Thus, the effective Lagrangian (3) can be identified with the amplitude of forward scattering of photons on atoms [57]

$$L^{\gamma a}_{\text{eff}} = N_{\text{at}} \sum_n \frac{M_n^a M_n^{\gamma}}{E_0 - E_n + \omega - i\Gamma_n/2},$$

where the summation goes over the atomic (or nuclear) excited states $n$ with the energies $E_n$ and widths $\Gamma_n$. 
Here $\omega$ is the energy of the incident photon, $M^a_{0n}$ is the matrix element of the axion-electron interaction \textsuperscript{[4]} and $M^a_{00}$ is the matrix element of the interaction of atomic electrons with external electromagnetic field of the incident photon.

The matrix element $M^a_{0n}$ in Eq. \textsuperscript{[5]} is given by $M^a_{0n} = \langle 0 | H^{ae} | n \rangle$, where the Hamiltonian $H^{ae}$ corresponds to the interaction \textsuperscript{[4]}. For non-relativistic electrons this Hamiltonian has the form \textsuperscript{[5],52}

$$H^{ae} = g_{ae} \left( \nabla \cdot \sigma - \partial_t \frac{p \cdot \sigma}{m_e} \right), \quad (6)$$

where $m_e$ is the electron mass and $p$ is its momentum. The form of this Hamiltonian can be further specified when the axion field is given by the plane wave with energy $\omega$ and wave-vector $k$, $a(r,t) = a_0 \sin(\omega t - kr)$. Then the Hamiltonian may be represented in the form of multipole expansion with the leading terms given by (modulo the oscillating factor $e^{i\omega t}$)

$$H^{ae} = g_{ae} a_0 \left( H^{ae}_{M0} + H^{ae}_{M1} + \ldots \right), \quad (7)$$

$$H^{ae}_{M0} = \frac{1}{2} \left( -\frac{\omega}{m_e} p \cdot \sigma + i(k \cdot r)(k \cdot \sigma) \right), \quad (8)$$

$$H^{ae}_{M1} = -\frac{1}{2}(k \cdot \sigma). \quad (9)$$

The scattering amplitude \textsuperscript{[5]} may be considered as a process, where an atom is (virtually) excited by absorbing a photon and returns back to the initial state via axion emission. Note that the M0 axions induce transitions between the states of opposite atomic parity but the same total angular momentum. Thus, M0 axions are produced when atoms absorb E1 photon which also changes the parity. Analogously, M1 axions do not change parity, and they may be produced when M1 photons are absorbed. As a result, in Eq. \textsuperscript{[5]} the matrix element of the operator $H^{ae}_{M0}$ should be coupled with the matrix element of E1 photon absorption operator $H^{ae}_{M1} = e (u \cdot r)$, while the matrix element of the operator $H^{ae}_{M1}$ should be multiplied by the matrix element of M1 photon-absorption operator $H^{ae}_{M1} = \frac{e}{2m_e} \langle n \times e \rangle (J + S)$, where $n$ is the propagation unit vector of the photon and $e$ is its polarization. The former case corresponds to effective magnetic field while the latter generates the effective electric field due to the interaction of photons with atoms,

$$B_{eff} = N_{at} \text{Re} \sum_n \frac{\langle 0 | H^{ae}_{M0} | n \rangle \langle n | H^{ae}_{M1} | 0 \rangle}{E_0 - E_n + \omega - i\Gamma_n/2}, \quad (10)$$

$$E_{eff} = N_{at} \text{Re} \sum_n \frac{\langle 0 | H^{ae}_{M0} | n \rangle \langle n | H^{ae}_{M1} | 0 \rangle}{E_0 - E_n + \omega - i\Gamma_n/2}. \quad (11)$$

A resonant atomic transition may lead to a significant photon absorption and a phase shift of the photon field (relative to the axions field) which destroys the coherence. Both effects may be seen from the complex phase in the photon field $A \propto e^{i\eta_n kr}$ where

$$n_r = 1 - 2\pi N_{at} \sum_n \frac{|M^\gamma_{0n}|^2}{E_0 - E_n + \omega - i\Gamma_n/2} \quad (12)$$

is the complex refractive index. One could go to the tail of resonance to suppress the photon absorption since the imaginary part of $n_r$ decreases quadratically with $E_0 - E_n + \omega$, but the phase decreases slowly (linearly), similar to $E_{eff}$ and $B_{eff}$ in Eqs. \textsuperscript{[10]} and \textsuperscript{[11]}.

### III. NUMERICAL ESTIMATES

Note that the effective fields \textsuperscript{[10]} and \textsuperscript{[11]} are proportional to the atom density $N_{at}$. Thus, a stronger effective field is produced in dense atomic systems such as solids, liquids or dense gases. In this note we consider only noble gases in liquid phase and heavy metal vapors since they can be treated using the standards atomic methods.

For this reason it is convenient to consider liquid xenon with atomic density $N_{at} = 1.3 \times 10^{22}$ cm$^{-3}$. The resonant M0 transition is possible from the ground state $0^+$ to the excited state $0^-$ with the energy $E_{0^-} = 9.45$ eV, while the corresponding photon transition is highly forbidden. This photon transition may be open when an external magnetic field $B_2$ is applied. This magnetic field causes the mixing of the states $0^-$ and $1^-$, where the latter corresponds to the energy level $E_{1^-} = 9.57$ eV. As we demonstrate in Appendix A the effective magnetic field may reach the value

$$B_{eff} \approx 1.2 \times 10^{-2} \text{ T}. \quad (13)$$

We stress that, according to Eq. \textsuperscript{[10]} the strength of interaction of this effective field with the axion field is specified by the coupling constant $g_{ae}$, which is independent from the axion-photon constant $g_{a\gamma}$. Thus, although this effective field is much weaker than the physical magnetic field $B = 5.3$ T applied at the ALPS experiment \textsuperscript{[58]}, it may be significant in the processes of axion production or detection if $g_{ae} \gg g_{a\gamma}$.

To estimate the effective electric field \textsuperscript{[11]} we consider a vapor of Tl atoms at temperature $T = 1200^\circ$C and density $N_{at} = 6.6 \times 10^{17}$ cm$^{-3}$. This system proved useful in the study of weak-interaction-induced optical activity of heavy-metal vapors \textsuperscript{[59]}. Indeed, the transition between the ground state $A$ with $J = \frac{1}{2}$ and the nearest excited state $B$ with $J = \frac{3}{2}$ is of M1 type. The energy difference between these states $\omega = 0.966$ eV corresponds to near-infrared region. In Appendix A we give an estimate of the effective electric field corresponding to such atomic transitions:

$$E_{eff} \approx 1.5 \times 10^5 \frac{V}{\text{cm}}. \quad (14)$$

### IV. CONCLUDING REMARKS

The aim of this note is to attract the attention to the possibility to have resonance enhancement of the axion-photon transformations in a medium. Such transitions between photons and axions in atoms can be represented...
via the effective electric and/or magnetic fields, similar to the interaction of the axion with physical electromagnetic fields. The crucial difference of these effective fields from the physical ones is that they couple to the axion through the independent coupling constant $g_{\mu \tau}$ which, for general axion-like particles, may be much larger than the axion-photon coupling $g_{\mu \tau}$. This motivates to set up new axion production/detection experiments where the medium plays the role of physical magnetic field in the current axion-search experiments. Such new experiments would allow one to measure the axion-electron coupling constant in the same way as the axion-photon coupling is studied. It is tempting to design such experiments and to estimate their efficiency.

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Appendix A: Estimates of effective fields

1. Effective magnetic field produced by liquid xenon

Let us consider the excited state $0^-$ in Xe atom with energy $E_{0^-} = \omega = 9.447$ eV. This state has zero total angular momentum, $J = 0$, and odd parity in contrast with the even-parity ground state $0^+$. Therefore, the axion transition between these states is of M0 type. The corresponding matrix element for resonant axion absorption was calculated in our recent work [55]

$$M^r = \langle 0^- | H_{a0}^{\mu r} | 0^+ \rangle = i \frac{\sqrt{2}}{3} \omega^2 R,$$  

(A1)

where $R$ is the radial integral $R = \int_0^{\infty} f_{0s_{1/2}}(r) f_{5p_{1/2}}(r) r^3 dr$. Here $f_{0s_{1/2}}(r)$ and $f_{5p_{1/2}}(r)$ are radial parts of $6s_{1/2}$ and $5p_{1/2}$ wavefunctions, respectively (see, e.g. [59]). For Xe atom, the value of this integral may be deduced from [60]. $R \approx -1.09$ a.u.

The photon transition $0^+ \leftrightarrow 0^-$ is highly forbidden. However, this transition may be opened by applying a magnetic field $B_2$ which provides the mixing of the excited states ($|0^+\rangle$ and $|1^+\rangle$),

$$|1\rangle = |0^+\rangle - i \frac{1}{\sqrt{2}} |\mu \cdot B_2 |0^+\rangle E_{0^-} - E_{1^-},$$  

(A2)

where $E_{1^-} = 9.57$ eV; $\mu = -\mu_B (J + S)$ is the electron magnetic moment operator and $\mu_B$ is the Bohr magneton. Taking into account that the matrix element $\langle 0^- | \mu \cdot B_2 | 1^- \rangle$ is $2/3 \mu_B B_2$, the photon E1 transition amplitude between the ground state $|0^+\rangle$ and the excited state $|1\rangle$ is found to be

$$M^\gamma = \frac{2}{3} \frac{e \mu_B B_2}{E_{0^-} - E_{1^-}} \langle 0^- | \epsilon \cdot r | 0^+ \rangle = \sqrt{2} \frac{e \mu_B B_2}{E_{0^-} - E_{1^-}} \frac{\epsilon \cdot \mu_B B_2 R}{9},$$  

(A3)

where $\epsilon$ is the photon polarization vector which is assumed to be along the $z$-axis, $\epsilon = (0, 0, 1)$. Thus, for the product of transition amplitudes (A1) and (A3) we find

$$M^a M^\gamma = \frac{2i \epsilon \mu_B B_2}{27 E_{0^-} - E_{1^-}} \omega^2 R^2.$$

Note that the expression (A4) is imaginary. Hence, the real part of the effective Lagrangian (5) corresponds to the resonant absorption. Note also that for liquid Xe the dominant contribution to the width of the state $\Gamma$ comes from the collisional broadening, $\Gamma \approx \Gamma_{\text{col}} = 2 \nu_0 N_{\text{at}} \sigma_{\text{col}}$, where $\nu_0 = \sqrt{2 k_B T/m_{\text{at}}}$ is the most probable thermal speed of the atoms and $\sigma_{\text{col}}$ is the collisional cross section. Here $k_B$ is the Boltzmann constant; $m_{\text{at}}$ and $R$ are atomic mass and radius, respectively. Assuming that the temperature of liquid Xe is $T = 164$ K and the atom density is $N_{\text{at}} = 1.3 \times 10^{22}$ cm$^{-3}$, the collisional width is estimated as $\Gamma_{\text{col}} \approx 1.4 \times 10^{-3}$ eV. Taking into account Eq. (A4), we find the expression for the effective magnetic field

$$B_{\text{eff}} = \frac{2}{27} N_{\text{at}} \frac{\mu_B B_2}{E_{0^-} - E_{1^-}} \frac{\omega^2 R^2}{\Gamma_{\text{col}}}.$$

Assuming that the magnetic field $B_2$ is of order of the magnetic field applied at ALPS experiment [58], $B_2 = 5$ T, we estimate the effective magnetic field produced by the atomic transitions (A5): $B_{\text{eff}} \approx 1.2 \times 10^{-2}$ T.

2. Effective electric field produced by thallium vapor

Consider the ground and excited states of the Tl atom, $A = |L = 1, J = \frac{3}{2}\rangle$ and $B = |L = 1, J = \frac{1}{2}\rangle$, with energy difference $\omega = 0.966$ eV. The transition between these states is of M1 type. The axion and photon absorption amplitudes for this transition are

$$M^a = -\frac{1}{2} \langle B| \kappa \cdot \sigma | A \rangle,$$

(A6)

$$M^\gamma = -\frac{\epsilon}{2 m_e} (n \times \epsilon) \langle B| J + S | A \rangle,$$

(A7)

where $n$ is the photon or axion propagation unit vector and $\epsilon$ is the photon polarization vector. For unpolarized atoms, in (11) we have to average over the states with different $z$-projection of the total angular momentum $M_J$. After this averaging, the contribution to the effective electric field vanishes, $\sum_{M_J} M^a M^\gamma = 0$, see [55].

To get a non-vanishing result in (11) one can apply a magnetic field $B_2$, which splits the Zeeman sublevels. The laser frequency is in resonance with a transition from only one of these sublevels. In particular, we suppose that the constant magnetic field $B_2$ is applied along the $z$-axis, $B_2 = B_2 \hat{z}$, and the laser frequency corresponds to the energy difference of levels $A = |1, \frac{3}{2}, -\frac{1}{2}\rangle$ and $B = |1, \frac{1}{2}, \frac{1}{2}\rangle$. We assume also that the photon propagates along the $y$-axis and is polarized in the $(x, z)$-plane.
In this case, the axion and photon matrix elements \( A_0 \) and \( A_1 \) are

\[
M_a = \frac{i}{3\sqrt{2}} \omega I, \quad (A8)
\]

\[
M^\gamma = \frac{\epsilon \omega I}{6\sqrt{2}m_e}, \quad (A9)
\]

where \( I \) is the overlap integral for 6p1/2 and 6p3/2 wavefunctions, \( I = \int_0^\infty dr r^2 f_{6p1/2}(r)f_{6p3/2}(r) \approx 0.98 \).

As is shown in Fig. 54, for T\(_1\) vapors at a temperature of \( T = 1473 \) K, the dominant contribution to the width of the line comes from Doppler broadening, \( \Gamma \approx \Gamma_{Dop} = 2v_0\omega/\sqrt{\pi} \), where \( v_0 = \sqrt{2k_BT/m_N} \) is the most probable thermal speed of the atoms. The numerical estimate for this width is \( \Gamma_{Dop} \approx 1.3 \times 10^{-6} \) eV. Taking this into account, we have the following expression for the effective electric field

\[
E_{\text{eff}} = N_{\text{at}} \frac{\epsilon \omega^2 \epsilon_5}{18m_e} \frac{\Gamma_{Dop}}{\Gamma_{Dop}}, \quad (A10)
\]

where the density of the vapor is \( N_{\text{at}} = 6.6 \times 10^{17} \) cm\(^{-3}\).

Assuming that the polarization of the photon is such that \( \epsilon_5 = 1 \), we estimate the effective electric field \( E_{\text{eff}} \approx 1.5 \times 10^5 \) V cm\(^{-1}\). Note that the same result may be obtained using the optical pumping rather than applying the magnetic field \( B_2 \).

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