# Field-Induced Axion Luminosity of Photon Gas via $a\gamma\gamma$-Interaction

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**Abstract**

The interaction of a pseudoscalar particle with two photons in an external electromagnetic field is used to study the photon decay $\gamma \to \gamma a$ and the photon coalescence $\gamma\gamma \to a$ where $a$ is a pseudoscalar particle associated with Peccei-Quinn $U(1)$ symmetry. A strong catalyzing influence of the external field on these processes reduces to that the field removes the suppression associated with the smallness of the axion mass. The field-induced axion emission by a photon gas is analyzed as one more possible source of energy losses by astrophysical objects.

## Introduction

Axion is a pseudoscalar particle which was introduced for solving the CP problem in QCD [1, 2]. This particle appears after the breakdown of the Peccei-Quinn chiral $U(1)$ symmetry and carries a small mass [3, 4]

$$10^{-5}\,\mathrm{eV} \lesssim m_a \lesssim 10^{-2}\,\mathrm{eV} \tag{1}$$

because the underlying symmetry is not exact at low energies. Now we can consider axions or any other pseudoscalar massless or low-mass bosons as a natural consequence of certain extensions of the standard model.

Axions couple to photons according to the Lagrangian [3]:

$$\mathcal{L} = \frac{g_{a\gamma}}{4}(F\tilde{F})a \tag{2}$$

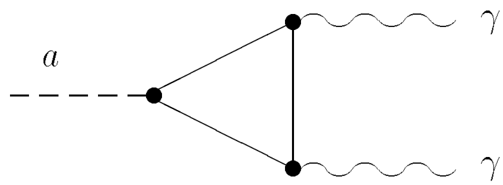


Figure 1: The effective vertex of the interaction of a pseudoscalar particle with two photons.

with a strength

$$g_{a\gamma} = \frac{\alpha}{\pi f_a} \xi, \tag{3}$$

where $f_a$ is the energy scale of the symmetry breaking and $\xi$ is a model-dependent factor of order unity. All existing axion models contain the interaction of an axion with charge fermions (usual or exotic) which automatically leads to an electromagnetic coupling of the form Eq. (2) because of the triangle fermion loop amplitude shown on Fig. 1.

The two-photon-axion interaction vertex allows for the axion radiative decay $a \to \gamma\gamma$ [3], for the Primakoff conversion $a \leftrightarrow \gamma$ in the presence of electric or magnetic fields [5] as well as for the photon decay $\gamma_T \to \gamma_L a$ [6, 7] and coalescence $\gamma_L \gamma_T \to a$ [7] in a plasma. The last two processes are kinematically possible because of the dispersion relations of electromagnetic excitations in a plasma which differ significantly from the vacuum dispersion [3].

Pseudoscalars are of great importance in an application to astrophysics as an additional source of star energy losses because of the very weak interaction with a matter. In astrophysical objects one has to take into account the influence of both components of the active medium, a plasma and an external electromagnetic field, on processes inside. A situation is also possible when the field component dominates and one can consider the axion processes in an external field only. Note that an arbitrary relatively smooth field in which a relativistic particle propagates is well described by the constant crossed field limit ($\mathbf{E} \perp \mathbf{B}$, $E = B$). In this case the dynamic parameter $\chi^2 = e^2(qFFq)/m^6$, where $m$ is a mass of an interactive particle (real or, possibly, virtual) with an electric charge $e$, is the only field invariant.

In a supernova explosion a region outside the neutrinosphere of order of hundred kilometers with a rather rarefied plasma with the temperature

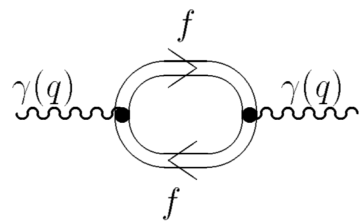


Figure 2: The diagram describing the polarization operator of a photon. Double lines correspond to fermion propagators in an external electromagnetic field.

of order of MeV and a strong magnetic field of order of the Schwinger value $B_e \simeq 4.41 \cdot 10^{13}$ G could exist. Under the supernova conditions considered the limit of small values of the dynamic parameter $\chi$ is realized in the processes of the axion emission by a photon gas.

In this talk we discuss the forbidden in vacuum axion processes – the photon decay $\gamma \to \gamma a$ and the photon coalescence $\gamma\gamma \to a$ in an external electromagnetic field using the effective $a\gamma\gamma$-interaction [8] and estimate their possible influence on the supernova cooling.

# Dispersions of Photon and Axion in External Electromagnetic Field

In calculation of probabilities and luminosities of the axion emission processes in an external field one has to take into account the non-trivial kinematics of interactive particles. The kinematics depends substantially on the particle dispersion relations which can change in the presence of the electromagnetic field because this field plays the role of an anisotropic medium. In this section the influence of the external electromagnetic crossed field on the photon [9] and axion [10] dispersions is considered.

The polarization operator of a photon in an external field can be obtained from the two-point fermion loop diagram shown on Fig. 2. The double solid lines denote the exact fermion propagators in an external electromagnetic field. The polarization operator can be presented as:

$$\Pi_{\alpha\beta} = i \sum_{\lambda=1}^{3} \Pi^{(\lambda)} \, \varepsilon_\alpha^{(\lambda)} \varepsilon_\beta^{(\lambda)}. \tag{4}$$

The set of the eigenvectors of $\Pi_{\alpha\beta}$ coincides with the photon polarizations:

$$\begin{aligned} \varepsilon_\alpha^{(1)} &= \frac{(qF)_\alpha}{\sqrt{(qFFq)}}, \quad \varepsilon_\alpha^{(2)} = \frac{(q\tilde{F})_\alpha}{\sqrt{(qFFq)}}, \\ \varepsilon_\alpha^{(3)} &= \frac{q^2 (qFF)_\alpha - (qFFq) q_\alpha}{\sqrt{q^2 (qFFq)^2}}. \end{aligned} \tag{5}$$

The first two vectors $\varepsilon^{(1)}$ and $\varepsilon^{(2)}$ describe the real transverse photon polarizations. In a difference with a plasma where besides the transverse photons the longitudinal excitation – "plasmon" – is appeared, in an external electromagnetic field the photon with the longitudinal polarization $\varepsilon^{(3)}$ is absent.

The eigenvalues $\Pi^{(\lambda)}$ with $\lambda = 1, 2$ of the polarization operator (4) determine the dispersion relations $q^2 - \Pi^{(\lambda)} = 0$ of the transverse photon excitations. The analysis of the polarization operator shows that $\Pi^{(\lambda)}$, in general, are complex $\Pi^{(\lambda)} = \mu_\lambda^2 - 2i\omega\Gamma_\lambda$. The real part $\mu_\lambda^2$ has a meaning of the field-induced "effective mass" squared and $\Gamma_\lambda$ is the probability of the photon decay $\gamma \to e^+e^-$.[1]

The photon dispersion curves in the external crossed field [9] are presented on Fig. 3 in dependence on the dynamic parameter $\chi$. The "effective masses" squared of the transverse photon polarization being alike in their qualitative behavior are different quantitatively. The difference in values of the "effective masses" squared of the photon eigenmodes makes possible the photon decay $\gamma \to \gamma a$ where $a$ is an arbitrary relatively light pseudoscalar. The analysis shows that in the physically interesting region of the dynamic parameter $\chi$ ($\chi \lesssim 10^2$) the photon "effective masses" squared are limited as:

$$|\mu_\lambda^2| \lesssim 10 \, \mathrm{keV}. \tag{6}$$

The field-induced contribution $\delta m_a$ to the small axion mass $m_a$ (1) can be calculated as the real part of $a \to e^+e^- \to a$ transition amplitude via

[1]We will consider the contribution of an electron only as the most sensitive to the external field fermion.

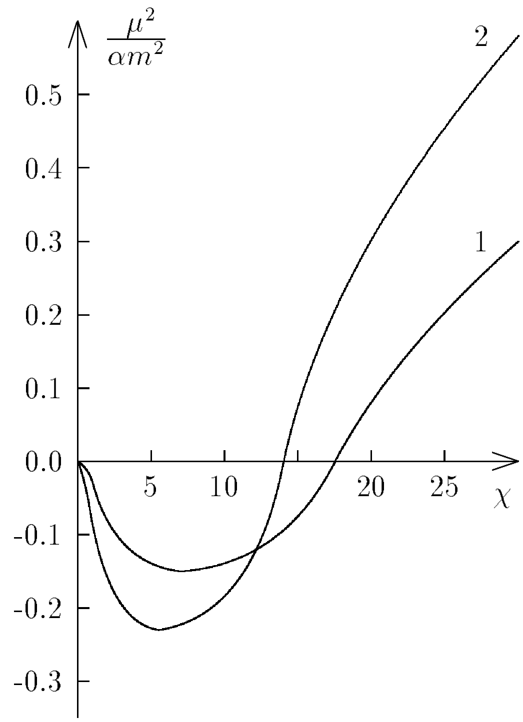


Figure 3: The photon dispersion curves in the crossed field.

electron loop. As a result $\delta m_a$ can be estimated as [10]:

$$\frac{\delta m_a^2}{m_a^2} \sim 10^{-10}\, C_e^2\, \chi^{2/3}, \tag{7}$$

where $C_e$ is the model-dependent factor which determines the electron-axion coupling $g_{ae}$. This contribution is negligibly small, and hereafter we will take the axion as a massless particle.

# Photon Decay $\gamma \to \gamma a$

The field-induced two-photon-axion vertex was obtained by us earlier [8] and used to analyze the axion radiative decay [11]. As in the case of $a \to \gamma\gamma$ the main contribution to the amplitude of the photon decay $\gamma \to \gamma a$ comes from the bilinear on the external field terms of $a\gamma\gamma$-vertex.

In the case of the small values of the dynamic parameters the decay of the photon with the first polarization $\varepsilon^{(1)}$ is allowed kinematically (see Fig. 3) due to the condition $\mu_1^2 > \mu_2^2$:

$$\gamma^{(1)} \to \gamma^{(2)} + a.$$

With the photon polarization vectors (5) the amplitude of the decay is:

$$M_d \simeq -\frac{4\alpha C_e m^2}{\pi f_a} t\,(1-t)\,\chi_1^2\, J(t\chi_1, \chi_1), \tag{8}$$
$$J(t\chi_1, \chi_1)\Big|_{\chi_1 \ll 1} \simeq \frac{2}{63}\,\chi_1^2\,(1-2t),$$

where $t = \omega_2/\omega_1$ is the relative energy of the final photon, $\chi_1$ is the dynamic parameter of the decaying photon. Taking all the particles as ultrarelativistic the probability of the photon decay is:

$$W_d^{(F)} = \frac{1}{16\pi\omega_1}\int\limits_0^1 dt\,|M_d|^2 \simeq 4.8\cdot 10^{-6}\,\frac{\alpha^2 C_e^2 m^4}{\pi^3 f_a^2 \omega_1}\,\chi_1^8. \tag{9}$$

The decay probability is proportional to the eighth power of the dynamic parameter $\chi_1$ and, hence, the eighth power of the external field strength.

# Photon Coalescence $\gamma\gamma \to a$

Both the photon decay $\gamma \to \gamma a$ and the photon coalescence $\gamma\gamma \to a$ give the contribution to the star energy losses due to the axion emission. The amplitude of the photon coalescence may be easily produced from the axion radiative decay amplitude [11] by changing all the 4-momenta of particles on opposite ones [12]:

$$M_c \simeq \frac{4\alpha C_e m^2}{\pi f_a}\,(\chi_1+\chi_2)\,[\chi_1 J(\chi_1,\chi_2) + \chi_2 J(\chi_2,\chi_1)], \tag{10}$$
$$\chi_1 J(\chi_1,\chi_2) + \chi_2 J(\chi_2,\chi_1)\Big|_{\chi_1,\chi_2 \ll 1} \simeq \frac{2}{21}\,\chi_1\chi_2(\chi_1+\chi_2),$$

where $\chi_1$ and $\chi_2$ are the dynamic parameters of the initial photons. Because the final state is the one-particle one the photon decay probability has the energy $\delta$-function and in the case of ultrarelativistic interactive particles is the following:

$$
\begin{aligned}
W_c^{(F)} &= \frac{2\pi\delta(\omega_1 + \omega_2 - E_a)}{8\omega_1\omega_2 E_a V} \, |M_c|^2 \\
&\simeq \frac{16\alpha^2 C_e^2 m^4}{441\pi f_a^2} \, \frac{\delta(\omega_1 + \omega_2 - E_a)}{\omega_1\omega_2 E_a V} \, \chi_1^2 \, \chi_2^2 \, (\chi_1 + \chi_2)^4,
\end{aligned} \tag{11}
$$

where $V$ is the normalized volume. The probability of the photon coalescence has the same dependence of the external field strength as the photon decay probability.

# Axion Luminosity of Photon Gas

To illustrate a possible astrophysical application of the results obtained we calculate the contributions of the photon decay and the photon coalescence to the axion emissivity $Q_a$ of a photon gas, i.e. the rate of energy losses per unit volume. The decay $Q_a^{(d)}$ [10] and coalescence $Q_a^{(c)}$ [12] axion emissivities are:

$$
\begin{aligned}
Q_a^{(d)} &= \int \frac{d^3 q_1}{(2\pi)^3} \, \omega_1 \, n_B(\omega_1) \int\limits_0^1 dt \, \frac{dW_d^{(F)}}{dt} \, (1 - t) \, (1 + n_B(\omega_1 t)) \qquad (12) \\
&\simeq 2.15 \, \frac{\alpha^2 C_e^2 m^7}{\pi^5 f_a^2} \left(\frac{T}{m}\right)^{11} \left(\frac{B}{B_e}\right)^8, \\
Q_a^{(c)} &= \frac{1}{2V} \int \frac{V d^3 q_1}{(2\pi)^3} \, n_B(\omega_1) \int \frac{V d^3 q_2}{(2\pi)^3} \, n_B(\omega_2) \, (\omega_1 + \omega_2) \, W_c^{(F)} \qquad (13) \\
&\simeq 121.85 \, \frac{\alpha^2 C_e^2 m^7}{\pi^5 f_a^2} \left(\frac{T}{m}\right)^{11} \left(\frac{B}{B_e}\right)^8,
\end{aligned}
$$

where $n_B(\omega_i)$ is the Planck distribution function of a photon with the energy $\omega_i$ at a temperature $T$, $B_e = 4.41 \cdot 10^{13}$ G is the Schwinger value. The factor $1/2$ in Eq. (13) takes into account the equivalence of the initial photons. The axion emissivity due to the effective $a\gamma\gamma$-vertex is determined by the photon coalescence because the contribution of the photon decay is the correction of order of a percent according to Eqs. (12) and (13).

The emissivities (12) and (13) allow to estimate the contribution of the considered processes into the axion luminosity $L_a$ (the energy losses by the escaping axions per unit time) in a supernova explosion from a region of order of hundred kilometers in size outside the neutrinosphere. In this region a rather rarefied plasma with the temperature of order of MeV and a magnetic field with the strength of order of $10^{13}$ G can exist. Under these conditions the estimation of the axion luminosity of a photon gas is:

$$L_a \simeq 2 \cdot 10^{38} \, \frac{\text{erg}}{\text{s}} \left(\frac{g_{ae}}{10^{-13}}\right)^2 \left(\frac{T}{1\text{MeV}}\right)^{11} \left(\frac{B}{10^{13}\,\text{G}}\right)^8 \left(\frac{R}{10^3\,\text{km}}\right)^3 . \qquad (14)$$

The comparison of Eq. (14) with the total neutrino luminosity from the neutrinosphere $L_\nu \sim 10^{52}$ erg/s shows that the contribution of the axion emission processes $\gamma \to \gamma a$ and $\gamma\gamma \to a$ by the photon gas to the energy losses in a supernova explosion is very small.

# Conclusions

In our talk the field-induced photon decay $\gamma \to \gamma a$ and coalescence $\gamma \to \gamma a$ are studied where $a$ is a light pseudoscalar particle. For the pseudoscalar particle we considered the most widely discussed particle, the axion, corresponding to the spontaneous breaking of the Peccei-Quinn symmetry. The forbidden in vacuum photon decay becomes kinematically possible because photons of different polarizations obtain different field-induced “effective masses” squared. At the same time an external field influence on the axion mass is negligible.

The processes $\gamma \to \gamma a$ and $\gamma\gamma \to a$ could be of interest as an additional source of energy losses by astrophysical objects. We considered the case of small values of the dynamic parameter which can be realized, for example, in a supernova explosion. The axion emission by the photon decay is the correction of order of a percent to the emission by the photon coalescence. While these processes and their evaluation are conceptually quite intruiging, the actual energy-loss rate appeared to be rather small in comparison with the neutrino luminosity in the conditions considered.

# Acknowledgements

The work of N.V. Mikheev was supported under grant No. d98-181 by International Soros Science Education Program. This work was partially supported by INTAS under grant No. 96-0659 and by the Russian Foundation for Basic Research under grant No. 98-02-16649.

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