Axion excitation by intense laser fields

To cite this article: J. T. Mendonça 2007 *EPL* **79** 21001

A LETTERS JOURNAL EXPLORING THE FRONTIERS OF PHYSICS

July 2007

EPL, **79** (2007) 21001 www.epljournal.org
doi: 10.1209/0295-5075/79/21001

Axion excitation by intense laser fields

J. T. MENDONÇA[(a)]

CFIF and GoLP, Instituto Superior Técnico - 1049-001 Lisboa, Portugal

received 14 May 2007; accepted in final form 9 June 2007
published online 11 July 2007

PACS 14.80.Mz – Axions and other Nambu-Goldstone bosons (Majorons, familons, etc.)
PACS 41.75.Jv – Laser-driven acceleration
PACS 95.35.+d – Dark matter (stellar, interstellar, galactic, and cosmological)

Abstract – Parametric excitation of axions by intense electromagnetic radiation, in the absence of any static field, is considered here. The possible use of this process, as the basis for a new concept of active experiments on axions, using the existing ultra-intense laser systems, is compared with the existing passive experiments using static magnetic fields.

Introduction. – The axions were proposed about twenty years ago to explain, in the frame of the standard model for particle physics, the observed charge-parity in strong interactions, or in other words, solve the strong *CP* problem [1–3]. These hypothetical elementary particles would have a very small mass (possibly in the meV range) and would couple very weakly with quarks, leptons and photons. For this reason they could be very good candidates for the explanation of the large amount of dark matter in the universe [4], and an important ingredient of cosmo-particle physics [5].

Various experiments are presently in operation, with the aim of discovering this new particle [6–8], using laboratory and astrophysical observations. Very recently a possible positive result was reported by the PVLAS experiment [9]. However, this result is dubious [10,11], and needs to be confirmed by an independent type of experiments, before concluding on the existence of this new particle [12]. Alternative explanations to the reported observations have been advanced, based on vacuum QED effects [13–16], but none of them is convincing or acceptable.

New results, based on independent experimental configurations are therefore necessary in order to arrive at a definite conclusion concerning the existence or not of axions. Experimental schemes based on X-ray laser facilities and large static magnetic fields were recently proposed [17]. Here we consider the direct axion-photon coupling in the absence of any static magnetic field, which could eventually lead to a new experimental concept. In contrast with the existing experimental arrangements using laser light [9,18], which considers rotating or modulated magnetic fields, this new concept would be based on the possible excitation of axions by two intense electromagnetic fields interacting in vacuum. We therefore explore a situation where the static rotating magnetic field of the PVLAS experiment is replaced by the electric and magnetic fields of an intense laser beam.

The interaction of axions with photons in the presence of a static magnetic field is theoretically well understood [19–21]. This corresponds to an indirect interaction, mediated by quarks, which is known as the Primakov effect. In this work we consider a generalized Primakov process, and establish the coupled mode equations describing the collective interactions between an axion and two photons, in the absence of any static field. Previous work in this subject was limited to massless pseudoscalar fields [22]. Explicit analytical results are derived, and used to estimate the possible observation of axions with the existing Peta-Watt lasers, or with the proposed Exa-Watt systems [23]. Comparison of the expected results with those recently reported in [9] will be made, and the experimental issues will be discussed.

Axion photon coupling. – Axions are elementary excitations of a pseudo-scalar field a. Photons couple to the axion field through the Lagrangian density [4]

$$L_{int} = \frac{1}{4} g_{a\gamma}\, F_{ij} \tilde{F}^{ij}\, a = g_{a\gamma} (\vec{E} \cdot \vec{B})\, a, \tag{1}$$

where $g_{a\gamma}$ is the coupling constant, F is the electromagnetic field tensor and $\tilde{F}$ its dual, with $\vec{E}$ and $\vec{B}$ the electric and magnetic fields. It is obvious that, under a *CP* transformation, this interaction Lagrangian will remain invariant because $\vec{E}$ is a polar vector and $\vec{B}$ is an axial one, therefore $(\vec{E} \cdot \vec{B})$ is a pseudo-scalar. Equation (1) means that an axion can couple to two photons, or in alternative

[(a)] E-mail: titomend@ist.utl.pt

to one photon in the presence of a static static electric or magnetic field. The first option will be considered here. This is described by coupled evolution equations that can be derived from the above Lagrangian. In terms of the electric and magnetic fields, the coupled equations are written (using $c=1$) as

$$\begin{aligned}\left(\partial_t^2-\nabla^2\right)\vec{E}&=-g_{a\gamma}\vec{B}\,\partial_t^2 a,\\ \left(\partial_t^2-\nabla^2+m_a^2\right)a&=g_{a\gamma}(\vec{E}\cdot\vec{B})\end{aligned}\tag{2}$$

where $\partial_t\equiv\partial/\partial t$ and m_a is the axion mass. It is obvious from these equations that two electromagnetic waves (or two different photon modes) with frequencies ω and ω' can only interact with the axion field if their polarization states are nearly orthogonal to each other, such that $\vec{E}(\omega)\cdot\vec{E}'(\omega')\simeq 0$. In this case they can exchange energy with the pseudo-scalar field a. Let us consider an axion field of the form

$$a(\vec{r},t)=a_0\exp(i\vec{k}_a\cdot\vec{r}-i\omega_a t)+\text{c.c.}\tag{3}$$

and two electromagnetic waves with field amplitudes $\vec{E}$ and $\vec{E}'$, such that the total electric field is

$$\vec{E}(\vec{r},t)=\vec{E}\exp(i\vec{k}\cdot\vec{r}-i\omega t)+\vec{E}'\exp(i\vec{k}'\cdot\vec{r}-i\omega' t)+\text{c.c.}\tag{4}$$

In the absence of coupling between fields ($g_{a\gamma}=0$), eqs. (2) lead to the dispersion relations for the three field modes, $\omega_a^2=k_a^2+m_a^2$, $k^2=\omega^2$, and $k'^2=\omega'^2$.

The amplitudes a_0, $\vec{E}$ and $\vec{E}'$ will be constant, and the associated magnetic field amplitudes will be given by $\vec{B}=(\vec{k}\times\vec{E})/\omega$ and $\vec{B}'=(\vec{k}'\times\vec{E}')/\omega'$. However, in the presence of coupling ($g_{a\gamma}\neq 0$), these amplitudes will become slowly varying functions of space and time. Using a perturbative approach, which is well justified due to the smallness of the coupling parameter $g_{a\gamma}$, we can assume that the above dispersion relations still hold. Following a well-known perturbation procedure [24], we can derive from the above wave equations for the fields, the corresponding evolution equations for the slowly varying amplitudes. These new equations take the form

$$\left(\partial_t+\frac{\vec{k}}{\omega}\cdot\nabla\right)\vec{E}=-ig_{a\gamma}\,\frac{\omega_a^2}{2\omega}\vec{B}'a_0\exp(i\phi)\tag{5}$$

with a similar equation for the field $\vec{E}'$, and

$$\left(\partial_t+\frac{\vec{k}_a}{\omega_a}\cdot\nabla\right)a=ig_{a\gamma}\,\frac{1}{2\omega_a}(\vec{E}\cdot\vec{B}'+\vec{E}'\cdot\vec{B})\exp(-i\phi).\tag{6}$$

The phase function ϕ is defined here as

$$\phi\equiv\phi(\vec{r},t)=\Delta\vec{k}\cdot\vec{r}-\Delta\omega t\tag{7}$$

with $\Delta\vec{k}=\vec{k}_a-(\vec{k}+\vec{k}')$ and $\Delta\omega=\omega_a-(\omega+\omega')$. It can easily be recognized that maximum nonlinear coupling between the axion and the photon fields occurs for $\Delta k=0$ and $\Delta\omega=0$, otherwise coupling will be damped by phase mixing. This exact phase matching condition $\phi=0$ corresponds to energy and momentum conservation in the axion-two photons decay interaction, as determined by $\vec{k}_a=\vec{k}+\vec{k}'$, and $\omega_a=\omega+\omega'$.

It is important to discuss the physical consequences of such stringent conditions. First, if the axion mass was negligible, $m_a=0$, we would be reduced to a purely one-dimensional problem. The two electromagnetic modes (for instance two laser beams) and the axion field that verify both conservation relations would propagate in the same direction. However, the existence of a finite axion mass implies that, given the condition $\Delta\omega=0$, a perfect phase matching such that $\Delta\vec{k}=0$ can only be achieved at an angle β between the two photon modes. This angle is determined by

$$\cos\beta=1-\frac{m_a^2}{2\omega\omega'}.\tag{8}$$

For a small mass such that $m_a\ll\omega_{1,2}$ this is a very small angle of the order of $\beta\simeq\pm m_a(\omega\omega')^{-1/2}$. In principle, we could also study the decay $\omega_a=\omega-\omega'$. But, in this case there is no real value for β satisfying the phase matching condition $\vec{k}_a=\vec{k}-\vec{k}'$. As shown by eq. (8), if we change the sign of ω', we will obtain $\cos\beta>1$. This is due to the lack of axion momentum, associated with the existence of a finite mass. Therefore, the case of an axion with energy smaller than that of any of the two photons is strictly forbidden for an infinite interacting distance between the field modes. But in practice, the interaction length is always finite. On the other hand, an oblique propagation is incompatible with large interaction distances, because the intersection between two laser beams is maximized for parallel propagation.

Axion excitation. – Let us then focus on plausible physical situations, where there is a perfect frequency matching between the three interacting field modes $\Delta\omega=0$, but not a complete phase matching. The one-dimensional configuration, with all field modes propagating along the same axis Oz is considered, because it maximizes the interaction length. This implies the existence of a small $\Delta\vec{k}$, which is of the order of the axion mass m_a. Let us focus on the sum frequency case, $\omega_a=\omega'+\omega$. Starting from (5)-(6), and assuming that the field amplitudes along the interaction region only depend on the space coordinate z, we obtain

$$\frac{\mathrm{d}}{\mathrm{d}z}E=ig_{a\gamma}\,\frac{\omega_a^2\sin\theta}{2\omega}a_0E'^*e^{i\Delta kz},\tag{9}$$

$$\frac{\mathrm{d}}{\mathrm{d}z}E'=ig_{a\gamma}\,\frac{\omega_a^2\sin\theta}{2\omega'}a_0E^*e^{i\Delta kz},\tag{10}$$

with $\Delta k=k_a-(k'+k)$, and

$$\frac{\mathrm{d}}{\mathrm{d}z}a_0=ig_{a\gamma}\,\frac{\sin\theta}{\omega_a}E\,E'e^{-i\Delta kz},\tag{11}$$

where θ is the angle between the electric fields of the two photon modes $\vec{E}$ and $\vec{E}'$. We have neglected group velocity dispersion, because the axion velocity (in units $c=1$) is nearly equal to one, $v_a = k/\omega \simeq 1 - m_a^2/2k^2$. This is valid for moderately large interaction distances L, such that $(k-k_j)L \ll 1$.

It is now useful to consider the case where one of the photon fields, say $\vec{E}'$ is associated with a very intense laser pulse, such that its amplitude can be taken as a constant during the interaction process. This is particularly valid for the present problem of an extremely week interaction between modes. It corresponds to the parametric approximation, well known from nonlinear optics. In this case, the above equations reduce to a simple system of two coupled equations

$$\frac{\mathrm{d}}{\mathrm{d}z}E = iw\, a_0\, \exp(i\Delta kz), \qquad \frac{\mathrm{d}}{\mathrm{d}z}a_0 = iw_a\, E\, \exp(-i\Delta kz), \tag{12}$$

with

$$w = g_{a\gamma}\,\frac{\sin\theta\omega_a^2}{2\omega}E'^*, \qquad w_a = g_{a\gamma}\,\frac{\sin\theta}{\omega_a}E'. \tag{13}$$

These coupled equations can easily be integrated. Assuming that the axion field amplitude is initially equal to zero, $a_0(z=0)=0$, we obtain the solutions

$$\begin{aligned} E(z) &= E(0)\cos(\Omega z)\exp(i\Delta kz/2),\\ a_0(z) &= iE(0)\frac{w_a}{\Omega}\sin(\Omega z)\exp(-i\Delta kz/2), \end{aligned} \tag{14}$$

with the nonlinear parameter defined by $\Omega^2 = w\,w_a + (\Delta k/4)^2$. This shows that, the stronger the pump intensity $|E'|^2$, the faster is the growth of the axion field. But the axion amplitude $a_0(z)$ also depends on the initial amplitude of the probe field $|E(0)|$. In contrast with the pump laser pulse, the probe laser pulse can have a much lower intensity. But the most interesting aspect of the above solutions (14) is that they show the decrease in the probe field amplitude, $E(z)$, and the increase of the axion field $a_0(z)$ at the expense of both the pump and the probe fields. This means that the excitation of axions could, in principle, be monitored by a slight decrease of the number of probe photons with frequency ω.

An experiment using intense laser pulses in the Peta-Watt domain, could then be envisaged. Due to the smallness of the expected coupling parameter $g_{a\gamma}$, the quantity Ωz will certainly be very small, and we can use the following approximation for the expected relative variation of the probe field:

$$\frac{\Delta E(z)}{E(0)} \simeq -\frac{1}{2}(\Omega z)^2, \tag{15}$$

where $\Delta E = E(z) - E(0)$. Assuming that $\omega_a \simeq 2\omega$ and using $\sin\theta \simeq 1$, we can write $\Omega^2 \simeq g_{a\gamma}^2 |E'|^2$. This is written in a unit system such that $\hbar = 1$ and $c = 1$. Turning now to the SI unit system, this can be explicitly written as

$$\Omega^2 = g_{a\gamma}^2 \hbar c^2 I_{pump}, \tag{16}$$

where I_{pump} is the pump laser intensity. The coupling constant is usually expressed in units of (GeV^{-1}), and the laser intensity in units of watt/cm^2, which leads to the final value of

$$\Omega\,[\mathrm{s}^{-1}] = 3\times 10^{-4}\sqrt{I_{pump}\,[\mathrm{watt/cm^2}]}\, g_{a\gamma}\,[\mathrm{GeV}^{-1}]. \tag{17}$$

Therefore, for the Peta-Watt laser systems presently in operation, such that $I_{pump} \simeq 10^{22}\,(\mathrm{watt/cm^2})$, and for an assumed value of $g_{a\gamma} \sim 3\times 10^{-6}\,(\mathrm{GeV}^{-1})$, we would obtain the value of $\Omega \simeq 10^2\,(\mathrm{s}^{-1})$, or $3\times 10^{-7}\,(\mathrm{m}^{-1})$. The use of the above process of axion excitation is therefore not unconceivable.

In terms of possible experimental signatures of the axion-photon coupling, it is also convenient to consider a second-order process, by which new photon states with frequency different from ω and ω' can be excited. This process corresponds to the nonlinear mixing of the pump laser field with the excited axion field, to give a secondary field E'' with frequency $\omega'' = 2\omega + \omega'$. This process can be described by an equation similar to (9). In the absence of phase mismatch, the amplitude of this secondary field is determined by

$$\frac{\mathrm{d}}{\mathrm{d}\tau}E'' = ig_{a\gamma}\,\frac{\omega_a^2 \sin\theta''}{2\omega''}a_0 E', \tag{18}$$

where $\sin\theta'' \simeq 1$, which means that the new field $\vec{E}''$ is parallel to the probe field $\vec{E}$. Replacing here the solution for the axion field as given by eq. (14), this can be easily integrated to give

$$E''(z) = g_{a\gamma}\,\frac{\omega_a^2}{2\omega''}\frac{E(0)}{w}\cos(\Omega z + \delta), \tag{19}$$

where the constant of integration δ is such that $E''(0)=0$. For very small values of $\Omega z \ll 1$, we obtain $|E''(z)/E(0)| \simeq (\Omega z)^2/2$. Comparison with eq. (15) shows that the amplitude of the new field is of the same order of the variation for the probe field amplitude $\Delta E(z)$. If we use a large probe field amplitude $E(0)$ with frequency ω', the possible observation of photons with a new frequency ω'' will be a favorable option for possible axion detection.

Let us now examine the difference frequency case $\omega_a = \omega' - \omega$. The mode coupled equations (12) are now replaced by

$$\frac{\mathrm{d}}{\mathrm{d}z}E = iw\, a_0^*\, \exp(i\Delta kz), \quad \frac{\mathrm{d}}{\mathrm{d}z}a_0^* = -iw_a E\, \exp(i\Delta kz). \tag{20}$$

with $\Delta k = k' - (k + k_a)$. The corresponding solutions for $a_0(0)=0$ are given by

$$E(z) = E(0)\left[\cosh(\Omega' z) - \frac{i\Delta k}{2\Omega'}\sinh(\Omega' z)\right]\exp(i\Delta k\, z/2) \tag{21}$$

and

$$a_0(z) = iE(0)\frac{w_a}{\Omega'}\sinh(\Omega z)\exp(-i\Delta k\, z/2), \tag{22}$$

with $\Omega'^2 = w\,w_a - (\Delta k/2)^2$. In contrast with the sum frequency, the frequency difference $\omega_a = \omega' - \omega$ is parametrically unstable, leading to an exponential growth of the excited axion field. However, for the plausible experimental situations where $\Omega' z \ll 1$, no significant difference with respect to the previous oscillating solutions is expected in terms of the final field amplitudes.

Sideband cascades. – In the above calculation we have assumed that the frequencies of the pump and probe laser fields, were of the same order. It is now interesting to consider a different case where the probe frequency (called now ω_0) is much larger than the pump one, ω'. This can be applied, for instance, to the plausible experimental configuration where the pump is an intense laser in the near infrared, and the probe is a high harmonic or an X-ray signal. In such a case we can give a more general description of the axion-photon coupling spectrum, showing the occurrence of sideband cascades. We consider possible axion and electromagnetic field spectra of the form

$$a(\vec{r},t) = \sum_n a_n \exp(ik_{an}\, z - i\omega_n t) + \text{c.c.}, \tag{23}$$

where n is an integer, and

$$\begin{aligned}\vec{E}(\vec{r},t) &= \vec{E}' \exp(ik'z - i\omega' t) \\ &\quad + \sum_n \vec{E}_n \exp(i\vec{k}_n \cdot \vec{r} - i\omega_n t) + \text{c.c.},\end{aligned} \tag{24}$$

where $\vec{E}'$ is the field of the intense pump laser as before, but $\vec{E}_0$ is now the probe field, and $\vec{E}_{n\neq 0}$ the possible secondary fields generated in the axion-photon coupling process. The above frequencies and wave numbers are defined by $\omega_n = \omega_0 + n\omega'$, $k_n = \omega_n$ and $k_{an} = \sqrt{\omega_n^2 - m_a^2} = k_n[1-(m_a/\omega_n)^2]^{1/2}$. Using the parametric approximation as before, and assuming $\vec{E}_0 \cdot \vec{E}' = \vec{E}_n \cdot \vec{E}' = 0$, in order to maximize the axion-photon coupling we arrive at the following coupled mode equations:

$$\begin{aligned}&\frac{\mathrm{d}}{\mathrm{d}z} E_n = -g_{a\gamma}\, \frac{i|E'|}{2\omega_n} \\ &\times \left[\omega_{n-1}^2 a_{n-1} e^{-i\Delta_{n-1} z + i\delta} + \omega_{n+1}^2 a_{n+1} e^{-i\Delta_{n+1} z - i\delta}\right]\end{aligned} \tag{25}$$

and

$$\frac{\mathrm{d}}{\mathrm{d}z} a_n = g_{a\gamma}\, \frac{i|E'|}{\omega_n} \left[E_{n-1} e^{i\delta} + E_{n+1} e^{-i\delta}\right] e^{i\Delta_n z}, \tag{26}$$

where $\Delta_n = m_a^2/2\omega_n$, and $E' = |E'| e^{i\delta}$. Now, assuming that $\omega' \ll \omega_0$, or $\omega_{n\pm1} \simeq \omega_n$, and introducing new field amplitudes such that

$$\begin{aligned}A_n(z) &= i^{-n} E_n(z) e^{-i\Delta_n z/2 - in\delta}, \\ \alpha_n(z) &= i^n \omega_n a_n(z) e^{-i\Delta_n z/2 - in\delta},\end{aligned} \tag{27}$$

we can get the simplified decoupled equations

$$\begin{aligned}\frac{\mathrm{d}^2}{\mathrm{d}\tau_n^2} A_n &= -(2A_n + A_{n+2} - A_{n-2}), \\ \frac{\mathrm{d}^2}{\mathrm{d}\tau_n^2} \alpha_n &= -(2\alpha_n + \alpha_{n+2} - \alpha_{n+2}),\end{aligned} \tag{28}$$

with $\tau_n = (g_{a\gamma}^2 |E'|^2 + \Delta_n^2)^{1/2} z/2 \sim g_{a\gamma} |E'| z/2$. This suggests the following approximate solutions for the axion and photon sideband fields, written in terms of Bessel functions as

$$E(z) \simeq E(0) \sum_{n=0} J_{2n}(\tau_{2n}(z))\, e^{i\Delta_{2n} z/2 + 2in(\delta - \pi/2)}, \tag{29}$$

and

$$a(z) \simeq \frac{2}{\omega_n} E(0) \sum_{n=0} J_{2n+1}(\tau_{2n+1}(z))\, e^{i\Delta_{2n} z/2 + i(2l+1)(\delta - \pi/2)}. \tag{30}$$

These solutions are valid for initial conditions such that only the probe field is present at $t=0$, $E_0(0) \neq 0$, and $E_{n\neq 0}(0) = a_n(0) = 0$. This shows that, if we start from a probe photon beam with frequency ω_0, we will generate odd axion field components a_{2n+1} and even sidebands of the photon field with amplitudes E_{2n}. Due to the smallness of $g_{a\gamma}^2 |E'|^2$, in a plausible experimental condition only the first two sidebands will eventually be seen. Nevertheless, identification of such a sideband spectrum would give a clear signature of the existence of any pseudo-scalar field.

Static magnetic field. – It is now useful to compare the efficiency of the above axion-photon parametric process with the passive detection of axions using a rotating static magnetic field, as explored in the recently reported experiments [9]. In this case the mode coupled equations will be similar to the previous ones, but with the pump laser fields E' and B' replaced by a rotating static magnetic field, $B_r = \text{const}$ and $E_r = 0$. Such a configuration was abundantly considered in the literature. Here we simply explore the analogies with the above pump laser case, and describe the axion and photon field spectra in more general terms than usually considered in the literature.

The relevant quantity in the case of a static rotating magnetic field is the projection of this field in the direction of polarization of the intense laser field, which is $B(t) = B_r \cos(\omega_r t)$ where $\omega_r \ll \omega_0$ is the rotation angular frequency. Going back to eqs. (2), and using eqs. (3)-(4) for the field spectra, we can derive along similar lines the following one-dimensional coupled mode equations:

$$\begin{aligned}\frac{\mathrm{d}E_n}{\mathrm{d}z} &= \frac{i}{4} g_{a\gamma}\, B_r \omega_n \left(a_{n-1} e^{i\phi_{n-1}} + a_{n+1} e^{i\phi_{n+1}}\right), \\ \frac{\mathrm{d}a_n}{\mathrm{d}z} &= i g_{a\gamma}\, \frac{B_r}{2\omega_n} \left(E_{n-1} e^{i\phi'_{n-1}} + E_{n+1} e^{i\phi'_{n+1}}\right),\end{aligned} \tag{31}$$

with the new phase factors $\phi_{n\pm1} = (k_{a(n\pm1)} - k_n) z$ and $\phi'_{n\pm1} = (k_{n\pm1} - k_{an}) z$. These phase factors are larger than the previous ones, because the spatially uniform magnetic field cannot provide momentum. But, for order of magnitude estimates and not too large interacting distances $z=t$, we can neglect the phase factors and use the following approximate solutions:

$$E_{n=2l}(z)=(-1)^n E_0(0) J_{2l}(\sqrt{2}\Omega_r z), \quad E_{2n+1}(z)=0, \tag{32}$$

and

$$a_{n=2l+1}(z)=-\frac{2i}{\omega_n} E_0(0) J_{2l+1}(\sqrt{2}\Omega_r z), \quad a_{2n}(z)=0, \tag{33}$$

where the coupling parameter is now determined by $\Omega_r = \sqrt{2} g_{a\gamma} B_r/2$. Here again, we expect a photon sideband cascade, with frequencies $\omega_{n=2l}=\omega_0+2l\omega_r$, and an axion cascade, with frequencies $\omega_{n=2l+1}=\omega_0+(2l+1)\omega_r$.

An estimate of the signal associated with the first sideband $\omega_2=\omega_0+2\omega_r$ can be obtained from the above solutions, for $\Omega_r\tau \ll 1$, as $E_2(z)\sim (g_{a\gamma}z)^2 B_r^2 E_0/4$, in agreement with previous estimates [19]. Comparing the relative amplitude variation $\delta_{static}=E_2/E_0$ with that due to the parametric axion excitation by strong laser fields, as determined by eqs. (15) and (20), $\delta_{laser}=\Delta E(z)/E(0)\simeq E''(z)/E(0)$. The result is

$$\delta_{static}/\delta_{laser} \simeq \left(\frac{B_r}{B'}\right)^2 \left(\frac{z_{static}}{z_{laser}}\right)^2, \tag{34}$$

where B' is the magnetic field associated with the intense pump laser. For an experiment such that report by [9], the magnetic field B_r is of the order of 5 tesla, while the magnetic field of a Peta-Watt laser can attain 10^5 tesla. This means that the parametric process is more efficient by a large factor than that of the static rotating field. However, the interaction length z_{static} is about 10^4 meters for the static rotating magnetic field case, where cavity resonators can be used. In contrast, Peta-Watt laser experiments are not compatible with the use of optical cavities in the interaction region, and possible the interaction lengths are at most of the order of 1 meter. This means that the combined balance between field strength and interaction distance, makes these two experimental concepts of nearly equal efficiency. A possible Peta-Watt experiment based on the concept of parametric excitation of axions by intense laser fields, could then be a good candidate for an alternative and independent assessment of the axion detection problem. If the signal observed by the reported PVLAS experiment is not a spurious signal but is due to the existence of a pseudo-scalar particle, then an experiment with no static magnetic field but based on Peta-Watt laser technology would be able to observe a similar (and eventually improved) signal and to establish a proof of principle for the existence of such a hypothetical particle.

Conclusions. – Active interaction of photons with axions, in the absence of any static magnetic field was considered here. The parametric coupling of an axion field with intense electromagnetic wave fields was shown to be possible. The case of a very strong electromagnetic wave, associated for instance to an intense laser beam, coupled to the axion field by a second laser beam was described. This process, involving a pump laser field, which can be considered as constant, and a probe laser field, is described by equations that are very similar to those used in nonlinear optics. The possible use of a very small angle between the two laser beams can also prevent dephasing due to the finite axion mass, at the cost of reducing the interaction length. Explicit analytical solutions were derived, and order of magnitude estimates were given.

The efficiency of this active approach to axions was compared with the currently used passive methods. In particular, static rotating magnetic fields are presently being used in axion motivated experiments. The advantages and disadvantages of the active method based on the use of intense Peta-Watt laser experiments were emphasized. The present work shows that a new experimental approach to the axion problem, with the same or better accuracy than the present experiments using static rotating magnetic fields, can be envisaged using Peta-Watt laser systems. This could eventually become a new area of research for the ultra-intense laser systems, and help to solve the ambiguities of the reported experimental findings.

REFERENCES

[1] Peccei R. D. and Quinn H. R., *Phys. Rev. Lett.*, **38** (1977) 1440.
[2] Weinberg S., *Phys. Rev. Lett.*, **40** (1978) 223.
[3] Wilczek F., *Phys. Rev. Lett.*, **40** (1978) 279.
[4] Raffelt G. G., *Stars as Laboratories for Fundamental Physics* (University of Chicago Press, Chicago) 1996.
[5] Khlopov M. Yu., Sakharov A. S. and Sudarikov A. L., *Nucl. Phys. B (Proc. Suppl.)*, **72** (1999) 105; Khlopov M. Yu., *Cosmoparticle Physics* (World Scientific, Singapore) 1999; Khlopov M. Yu. and Rubin S. G., *Cosmological Pattern of Microphysics in Inflationary Universe* (Kluwer Academic Publishers, Dordrecht) 2004.
[6] van Bibber K. and Rosenberg L. J., *Phys. Today*, **59** August issue (2006).
[7] Zioutas K. *et al.*, *Phys. Rev. Lett.*, **94** (2005) 121301.
[8] Fairbairn M., Rashbar T. and Troitsky S., arXiv:astro-ph/0610844v1 (2006).
[9] Zavattini E. *et al.*, *Phys. Rev. Lett.*, **96** (2006) 110406.
[10] Melissinos A. C., arXiv:hep-ph/0702135v1 (2007).
[11] Mohapatra R. N. and Nasri S., *Phys. Rev. Lett.*, **98** (2007) 050402.
[12] Lamoreaux S., *Nature*, **441** (2006) 31.
[13] Mendonça J. T., Dias de Deus J. and Ferreira P. C., *Phys. Rev. Lett.*, **96** (2006) 100403; **97** (2006) 269901.
[14] Adler S. L., arXiv:hep-ph/0611267 (2006).
[15] Mendonça J. T., in preparation.
[16] Gies H., Jaeckel J. and Ringwald A., *Phys. Rev. Lett.*, **97** (2006) 140402.

[17] Rabadan R., Ringwald A. and Sigurdson K., *Phys. Rev. Lett.*, **96** (2006) 110407.
[18] Cameron R. *et al.*, *Phys. Rev. D*, **47** (1993) 3707.
[19] Raffelt G. and Stodolski L., *Phys. Rev. D*, **37** (1988) 1237.
[20] Maiani L., Petronzio R. and Zavattini E., *Phys. Lett. B*, **175** (1986) 359.
[21] Gasperini M., *Phys. Rev. Lett.*, **59** (1987) 396.
[22] De Sabbata V. and Gasperini M., *Lett. Nuovo Cimento*, **30** (1981) 363.
[23] Mourou G. A., Tajima T. and Bulanov S. V., *Rev. Mod. Phys.*, **78** (2006) 309.
[24] Boyd R. W., *Nonlinear Optics* (Academic Press, San Diego) 1992.