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Extending the reach of axion-photon regeneration experiments towards larger masses with phase shift plates

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Abstract

We present a scheme to extend the sensitivity of axion-photon regeneration experiments towards larger masses with the help of properly chosen and placed phase shift plates.

Many proposals to embedd the standard model of particle physics into a more general, unified framework predict a number of new very light particles which are very weakly coupled to ordinary matter. Typically, such light particles arise if there is a global continuous symmetry that is spontaneously broken in the vacuum – a notable example being the axion [1,2], a pseudoscalar particle arising from the breaking of a U(1) Peccei-Quinn symmetry [3] introduced to explain the absence of CP violation in strong interactions. Other examples of light spin-zero bosons beyond the standard model are familons [4], Majorons [5,6], the dilaton, and moduli, to name just a few. We will call them axion-like particles, ALPs, in the following.

At low energies, the coupling of such an ALP, whose corresponding quantum field we denote by ϕ, to photons is described by an effective Lagrangian,

$$\mathcal{L} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{1}{2}m_\phi^2\phi^2 - \frac{1}{4}g\phi F_{\mu\nu}\tilde{F}^{\mu\nu}\,, \tag{1}$$

where $F_{\mu\nu}$ ($\tilde{F}_{\mu\nu}$) is the (dual) electromagnetic field strength tensor[1] and m_ϕ is the mass of the ALP. Correspondingly, in the presence of an external magnetic field, a photon of energy ω may oscillate into an ALP of small mass $m_\phi < \omega$, and vice versa [9,10].

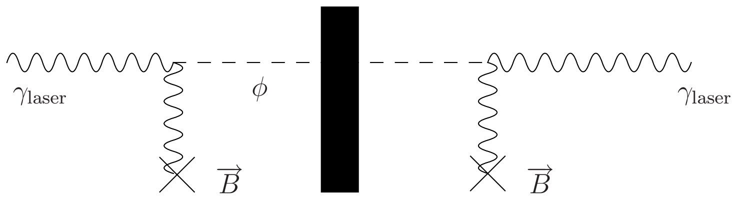

Figure 1: Schematic view of ALP production through photon conversion in a magnetic field (left), subsequent travel through an optical barrier, and final detection through photon regeneration (right).

The exploitation of this mechanism is the basic idea behind ALP-photon regeneration – sometimes also called "light shining through a wall" – experiments [11–13] (cf. Fig. 1). Namely, if a beam of photons is shone across a magnetic field, a fraction of these photons will turn into ALPs. This ALP beam could then propagate freely through an optical barrier without being absorbed, and finally another magnetic field located on the other side of the wall could transform some of these ALPs into photons – apparently regenerating these photons out of nothing.

A pioneering experiment of this type was carried out in Brookhaven by the Brookhaven-Fermilab-Rochester-Trieste (BFRT) collaboration, using two prototype magnets for the Colliding Beam Accelerator [14,15]. Presently, there are worldwide several second generation ALP-photon regeneration experiments under construction or serious consideration (cf. Table 1; for a review, see Refs. [21,22]). These efforts are partially motivated by

[1] The effective Lagrangian (1) applies for a pseudoscalar ALP, i.e. a spin-zero boson with negative parity. In the case of a scalar ALP, the $F_{\mu\nu}\tilde{F}^{\mu\nu}$ in Eq. (1) is replaced by $F_{\mu\nu}F^{\mu\nu}$. For the more general case where ϕ ceases to be an eigenstate of parity [7], see Ref. [8].

Table 1: Experimental parameters of upcoming photon regeneration experiments: magnetic fields B_i and their length ℓ_i on production ($i = 1$) and regeneration ($i = 2$) side (cf. Fig. 1).

Name	Laboratory	Magnets	Laser
ALPS [16]	DESY/D	$B_1 = B_2 = 5$ T $\ell_1 = \ell_2 = 4.21$ m	$\omega = 2.34$ eV
BMV [17]	LULI/F	$B_1 = B_2 = 11$ T $\ell_1 = \ell_2 = 0.25$ m	$\omega = 1.17$ eV
LIPSS [18]	Jlab/USA	$B_1 = B_2 = 1.7$ T $\ell_1 = \ell_2 = 1$ m	$\omega = 1.17$ eV
OSQAR [19]	CERN/CH	$B_1 = B_2 = 11$ T $\ell_1 = \ell_2 = 7$ m	$\omega = 1.17$ eV
PVLAS [20]	Legnaro/I	$B_1 = 5$ T $\ell_1 = 1$ m $B_2 = 2.2$ T $\ell_2 = 0.5$ m	$\omega = 1.17$ eV

the report from the PVLAS collaboration of evidence for a non-zero apparent rotation of the polarization plane of a laser beam after passage through a magnetic field [23]. While the size of the observed effect greatly exceeds the expectations from quantum electrodynamics [24–26], it is compatible with the expectations [27] arising in the context of a photon-ALP oscillation hypothesis. Indeed, the rotation observed by PVLAS can be reconciled with the non-observation of a signal by BFRT, if there exists an ALP with a mass $m_\phi \sim$ meV and a coupling $g \sim 10^{-6}$ GeV^{-1} [28]. Although these parameter values seem to be in serious conflict with bounds coming from astrophysical considerations, there are various ways to evade them [29–35]. Therefore, it is extremely important to check the ALP interpretation of PVLAS by purely laboratory experiments [32]. Moreover, it would be nice if in this way one might ultimately extend the laboratory search for ALPS to previously unexplored parameter values (see also Ref. [42]). In this letter, we propose a method to extend the sensitivity of the planned photon-regeneration experiments to higher ALP masses.

Let us start with an outline of the calculation of the photon $\to$ ALP conversion probability $P_{\gamma\to\phi}$, to lowest order in the coupling g. As emphasized in Ref. [13], this calculation amounts to solving the classical field equations following from Eq. (1),

$$\partial_\mu F^{\mu\nu} = g\partial_\mu \left(\phi\tilde{F}^{\mu\nu}\right); \qquad \left(\partial_\mu\partial^\mu + m_\phi^2\right)\phi = g\vec{E}\cdot\vec{B}\,, \tag{2}$$

to lowest order in $gB\ell$, where ℓ is the linear dimension associated with the extent of the magnetic field[2]. This can be done by neglecting the modification of the electromagnetic field due to the presence of the pseudoscalar field (through the right hand side of the first

[2]In the case of a scalar ALP, the term $\vec{E}\cdot\vec{B}$ in Eqs. (2), (3), and (4) is replaced by $\frac{1}{2}(\vec{E}^2-\vec{B}^2)$.

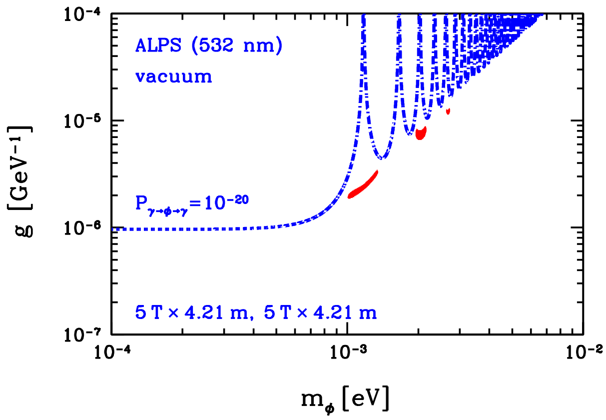

Figure 2: Two photon coupling g of the (pseudo-)scalar versus its mass m_ϕ. Iso-contour of the regeneration probability $P_{\gamma\to\phi\to\gamma} = P_{\gamma\to\phi}P_{\phi\to\gamma}$, for the parameters of the ALPS experiment, i.e. magnetic fields $B_1 = B_2 = 5$ T, over a length $\ell_1 = \ell_2 = 4.21$ m, exploiting a green ($\lambda = 532$ nm) photon beam, corresponding to $\omega = 2.34$ eV, in vacuum. Also shown in red are the 5 sigma allowed regions [28] from PVLAS data on rotation [23] plus BFRT data on rotation, ellipticity, and regeneration [15] plus Q&A data on rotation [38].

equation above). Solving for ϕ in the second equation yields [9, 13]

$$\phi^{(\pm)}(\vec{x},t) = \mathrm{e}^{-\mathrm{i}\omega t}\int \mathrm{d}^3x' \,\frac{1}{4\pi}\,\frac{\mathrm{e}^{\pm\mathrm{i}k_\phi|\vec{x}-\vec{x}'|}}{|\vec{x}-\vec{x}'|}\,g\vec{E}(\vec{x}')\cdot\vec{B}(\vec{x}')\,, \tag{3}$$

where the energy ω and the modulus of the three-momentum k_ϕ are related by $k_\phi = \sqrt{\omega^2 - m_\phi^2}$. This solution simplifies even more if we specialize to the usual experimental configuration of a laser photon beam send along the x-axis with fixed linear polarization in the z direction. If the transverse extent of the magnetic field is much larger than that of the laser beam, the problem is effectively one-dimensional. In one dimension and taking into account only ALPs that propagate into the positive x-direction, Eq. (3) becomes,

$$\phi^{(+)}(x,t) = \mathrm{e}^{-\mathrm{i}(\omega t - k_\phi x)}\,\frac{\mathrm{i}g}{2k_\phi}\int \mathrm{d}x'\,\vec{E}(x')\cdot\vec{B}(x'). \tag{4}$$

Inserting in Eq. (4) furthermore the appropriate plane wave form $\vec{E}_0(\vec{x},t) = \vec{e}_z E_0 \mathrm{e}^{\mathrm{i}\omega(x-t)}$ for the electric field of the laser beam and assuming, as realized in all the proposed experiments, a magnetic field with fixed direction along the z-axis and possibly variable (as a function of x) magnitude, $\vec{B}_0(\vec{x}) = \vec{e}_z B_0(x)$, one ends up with the solution[3]

$$\phi^{(\pm)}(\vec{x},t) = \frac{\mathrm{i}g}{2k_\phi}E_0\,\mathrm{e}^{-\mathrm{i}(\omega t - k_\phi x)}\int \mathrm{d}x'\,\mathrm{e}^{\mathrm{i}qx'}B_0(x')\,, \tag{5}$$

[3]The solution (5) applies also in the case of a scalar ALP, if the magnetic field direction is chosen to point into the y direction, $\vec{B}_0(\vec{x}) = \vec{e}_y B_0(x)$ (or, alternatively, if the polarization of the laser is chosen to point in the y direction).

where

$$q = k_\gamma - k_\phi = \omega - \sqrt{\omega^2 - m_\phi^2} \approx \frac{m_\phi^2}{2\omega} \tag{6}$$

is the momentum transfer to the magnetic field, i.e. the modulus of the momentum difference between the photon and the ALP. The probability that a photon converts into an axion-like particle and vice versa can be read off from Eq. (5) and reads [9,13]

$$P_{\gamma\to\phi} = P_{\phi\to\gamma} = \frac{1}{4}\,\frac{\omega}{k_\phi}\,g^2 \left| \int \mathrm{d}x'\, \mathrm{e}^{\mathrm{i}qx'}\, B_0(x') \right|^2 , \tag{7}$$

which reduces, for a constant magnetic field, $B_0(x') = \text{const}$, of linear extension ℓ, to

$$P_{\gamma\to\phi} \approx g^2\, B_0^2\, \sin^2\left(q\ell/2\right)/q^2\,. \tag{8}$$

Clearly, in the experimental setup considered, the maximum conversion probability, $P_{\gamma\to\phi} \approx g^2 B_0^2 \ell^2$, is attained at small momentum transfer, $q = m_\phi^2/(2\omega) \ll 1$, corresponding to a small ALP mass. For this mass range, the best limits are obtained in a straightforward manner by exploiting strong and long dipole magnets, as they are used for storage rings such as HERA [36] or LHC [37], cf. the experiments ALPS [16] and OSQAR [19], respectively (see Table 1). However, for larger masses, the sensitivity of this setup rapidly diminishes.

We illustrate this in Fig. 2, which displays an iso-contour of the light shining through a wall probability in the g-m_ϕ plane, exploiting the experimental parameters of the ALPS experiment [16]. Clearly, for this setup, the parameter region in g vs. m_ϕ suggested by the combination of BFRT plus Q&A exclusion and PVLAS evidence can not be probed. This is even more dramatic for the OSQAR experiment, which exploits an LHC magnet. Moreover, increasing the refraction index by filling in buffer gas does not help since it works in the wrong direction (contrary to the claim[4] in the ALPS letter of intent [16]).

A simple possibility to probe the meV region[5] in the ALPS setup is to reduce the effective length of the magnetic field region both on the production and detection side of the magnet by shortening the beam pipe on both sides. As can be seen in Fig. 3, this possibility enables to extend the mass region probed by the experiment, however at the expense of sensitivity: one looses about one order of magnitude in the light shining through a wall probability.

Another idea to extend the sensitivity towards larger ALP masses was introduced in Ref. [13]. There, it was shown that a segmentation of the magnetic field into regions

[4]In a refractive medium, the laser beam has a phase velocity $1/n \equiv v \equiv \omega/k_\gamma$. The momentum transfer (6) reads then $q = n\,\omega - \sqrt{\omega^2 - m_\phi^2} \approx \frac{m_\phi^2}{2\omega} + (n-1)\,\omega$. The second term in this expression has the opposite sign as the corresponding term in Ref. [16]. Correspondingly, one would need a buffer gas with refraction index less than unity, i.e. a plasma, in order to decrease q (and thereby maximize the conversion probability (8)) rather than to increase it. A.R. would like to thank Aaron Chou for pointing out the correct sign.

[5]Another possibility to probe larger ALP masses even with a long magnet would be to exploit VUV or X-ray free-electron laser beams [39–41]. However, at the moment conventional lasers seem to offer better prospects (see also Ref. [42])

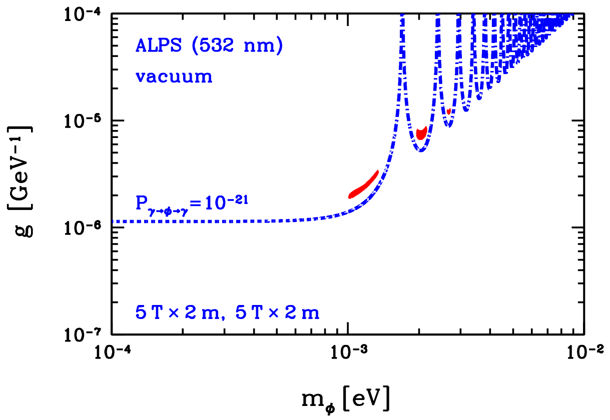

Figure 3: Iso-contour of the regeneration probability, as in Fig. 2, but with reduced lengths of the magnetic field region. Note, that the regeneration probability is reduced by a factor of 10.

of alternating polarity gives a form factor $\int \mathrm{d}x' \exp(\mathrm{i}qx')B_0(x')$ that peaks at a nonzero value of q, thereby giving sensitivity to higher-mass pseudoscalars. In fact, the conversion probability (7) reads [13, 43], in a magnet with N segments of alternating field direction (but the same magnitude B_0),

$$P_{\gamma\to\phi} \approx g^2 B_0^2 \frac{\sin^2(qd/2)}{q^2} \left| \sum_{k=1}^{N} (-1)^k \exp\{\mathrm{i}(2k-1)qd/2\} \right|^2 \tag{9}$$

$$= \frac{g^2 B_0^2}{q^2} \begin{cases} \sin^2(q\ell/2)\tan^2(q\ell/(2N)) & \text{for} \quad N \text{ even} \\ \cos^2(q\ell/2)\tan^2(q\ell/(2N)) & \text{for} \quad N \text{ odd} \end{cases},$$

where $d = \ell/N$ is the length of each of the N segments. For $N > 1$, this indeed gives rise to more sensitivity at non-zero values of q.

In this letter, we will introduce a similar, but more practical possibility based on the use of phase shift plates. The idea is very simple. From our starting point, Eq. (4), we can see that what counts is actually $\vec{E}(x')\cdot\vec{B}(x')$. The configuration based on N alternating magnetic fields is therefore equivalent to a configuration with non-alternating magnetic field, however with $N-1$ retardation plates with phase shift π ("$\lambda/2$" plates) inverting the sign of the electric field, placed equidistantly over the length ℓ of the magnet. In this case we have alternating signs of $\cos\theta$, where θ is the angle between $\vec{E}$ and $\vec{B}$, instead of alternating signs of the magnetic field. But both cases have an identical profile of $\vec{E}(x')\cdot\vec{B}(x')$. In Fig. 4, we show that with a proper choice of the number and positions of such phase shifters, ALPS should easily cover the region of parameter space suggested by PVLAS + BFRT + Q&A. The same applies for OSQAR.

Let us now get a more intuitive understanding of how this works and see how we can do even better. The crucial part in Eq. (5) is the integral

$$f(q) = \int \mathrm{d}x' \mathrm{e}^{\mathrm{i}qx'} B_0(x'). \tag{10}$$

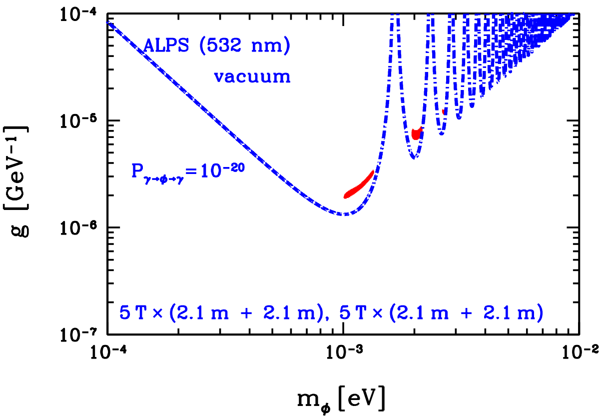

Figure 4: Iso-contour of the regeneration probability, as in Fig. 2. Here, we used one phase shift ("$\lambda/2$") plate each in the middle of the generation and the regeneration sides.

For a constant magnetic field of length ℓ the oscillating factor $\mathrm{e}^{\mathrm{i}qx'}$ suppresses the integral compared to the massless case with $q = 0$, where the integral is simply

$$|f(q)| < |f(0)| = \left|\int_0^\ell \mathrm{d}x' B_0\right| = B_0\ell, \quad \text{for } B_0(x) = \text{constant}. \tag{11}$$

This suppression arises because coherent production of ALPs works only if the ALP and the photon are in phase. The factor $\mathrm{e}^{\mathrm{i}qx'}$ accounts for the phase difference between ALP and photon.

To improve the situation one would want to bring photon and ALP back into phase with each other. This can be achieved in a simple way by the introduction of phase shift plates. A simplified picture of the phase correction process is given in Fig. 5. At the beginning, photon (red) and ALP (black) are in phase. However, due to its mass, the ALP has a slightly larger wavelength than the photon. After a few oscillations photon and ALP are more and more out of phase. Then we insert the phase shift plate (turquoise). With refractive indices $n > 1$ we cannot make the wavelength larger. So it is not possible to "delay" the photon until the ALP has caught up. What we can do, however, is to is to *increase* the phase difference between ALP and photon such that it is exactly 2π (the photon does an extra wiggle in Fig. 5.). Now, a phase shift of 2π is exactly equivalent to a phase shift of 0. Photon and ALP are in phase again. Therefore, we can keep photon and ALP in phase over quite long distances simply by inserting a suitable phase shift plate whenever the phase difference becomes too large, and we get coherent production over the whole length of the magnet.

Let us now understand more quantitatively how this works. To derive Eqs. (5) and (10) we have assumed that the photon is a plane wave. Therefore, we can identify $qx' = (k_\gamma - k_\phi)x'$ as the phase difference between the photon wave and the ALP wave at the point x'. In general we should write (cf. Eq. (4))

$$f(q) = \int \mathrm{d}x' \mathrm{e}^{i(\varphi_\gamma(x') - \varphi_\phi(x'))} B_0(x'), \tag{12}$$

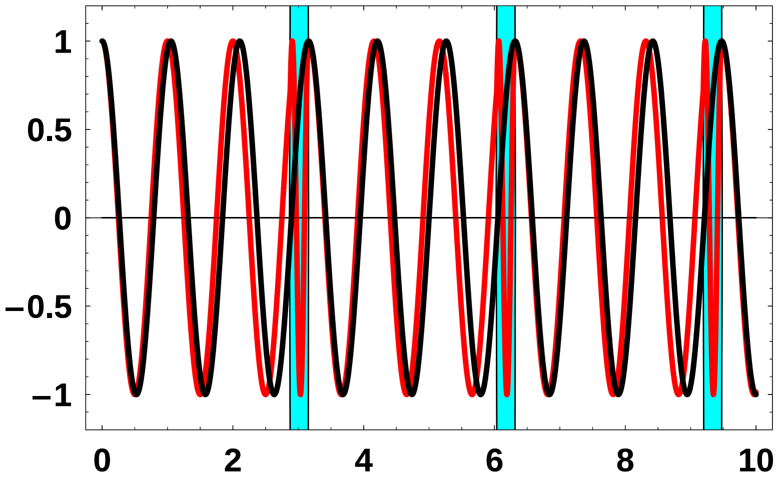

Figure 5: Illustration of the effect of a properly chosen and placed phase shift plate on the phase relation between photon and ALP (this simplified picture shows only the phase relation; the amplitudes of photon and ALP are not correct in this picture). Photon (red) and ALP (black) start in phase. Due to their different wavelength they are, however, somewhat out of phase after several oscillations - say by an amount ζ. This is corrected by introduction of a phase shift plate that causes the photon to get an extra phase $2\pi-\zeta$. In other words the plate causes the photon to complete the extra wiggle.

where $\varphi_{\gamma,\phi}$ are the phases of the photon and the ALP fields, respectively.

Let us imagine a situation where we insert $N-1$ thin[6], non-reflective[7] plates that accelerate the photon phase by κ at equidistant places $s\ell/N = s\Delta x$, $s = 1\dots N-1$, in a constant magnetic field of length ℓ. The plates affect only the photon. The ALP phase remains unaffected. Therefore, we have,

$$\begin{aligned} \varphi_\gamma(x) &= k_\gamma x + s\kappa \quad \text{for} \quad s\Delta x < x \leq (s+1)\Delta x, \qquad \Delta x = \frac{\ell}{N}, \\ \varphi_\phi(x) &= k_\phi x. \end{aligned} \tag{13}$$

Inserting this into Eq. (12), we find

$$f(q) = B_0 \sum_{s-0}^{N-1} \int_{s\Delta x}^{(s+1)\Delta x} \mathrm{d}x' \mathrm{e}^{\mathrm{i}(qx'+s\kappa)} = B_0 \mathrm{e}^{\frac{\mathrm{i}}{2}(q\ell+(N-1)\alpha)} \frac{2\mathrm{i}\sin\left(\frac{q\Delta x}{2}\right)}{q} \frac{\sin\left(\frac{N}{2}(q\Delta x+\kappa)\right)}{\sin\left(\frac{1}{2}(q\Delta x+\kappa)\right)} \tag{14}$$

[6]One might ask what happens to the photon-ALP system inside the material of the plate. One can check (cf., e.g., Ref. [10]) that for sufficiently large refractive index of the material, $n-1 \gg m_\phi^2/(2\omega^2)$, the mixing between photon and ALP is effectively switched off compared to the mixing in vacuum. Photon and ALP simply propagate through the plate without changing their amplitudes (the phases change, of course). In other words, the thickness of the plates has to be subtracted from the total length of the production or regeneration region. That is why we require thin plates. For practical purposes, this is a rather mild constraint. For $n-1 \sim 0.1$, the thickness of the plates required for a phase shift of the order of 2π is only $d \sim 10\,\lambda \sim 10\,\mu m$, which is tiny compared to the typical lengths of the production/regeneration regions which are of the order of a few m.

[7]Reflected photons are effectively lost.

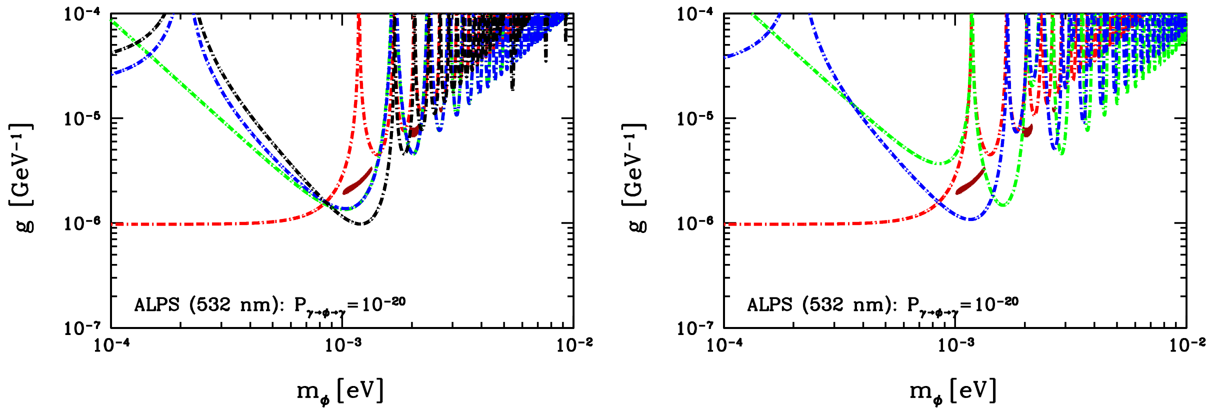

Figure 6: Iso-contours of the regeneration probability, as in Fig. 2. In the left figure, we have used no phase correction (red), one plate with $\kappa = \pi$ (green), and one plate with the optimal choice of κ according to Eq. (16) for $m_\phi = 1.2\,\mathrm{meV}$ (blue). The black curve is for 20 plates with the optimal choice of κ. In the right figure, we have the same but with 3 plates for the green and blue curves.

We can now choose the number of plates N and the phase shift κ according to the recipe described above. First we choose N large enough such that

$$\frac{1}{2}q\Delta x \ll 1 \tag{15}$$

And then we choose κ such that the phase difference that has accumulated over Δx is "completed" to 2π,

$$\kappa = 2\pi - \frac{1}{2}q\Delta x. \tag{16}$$

Evaluating Eq. (14) in the limit $N(q\Delta x + \kappa)/2 \to 0$ one finds

$$|f(q)| = B_0 \Delta x\, N = B_0 \ell. \tag{17}$$

And we have coherent production over the whole length ℓ.

The potential of this approach is demonstrated in Fig. 6 for the example of the ALPS experiment. In the optimized mass region we get more than ten times as many regenerated photons as we would get if the length of the magnet is reduced as in Fig. 3.

Another practical advantage of this method is that we can scan through a whole mass range. Performing several measurements with different phase shift plates we can always choose for each q, i.e. for each m_ϕ, plates with an appropriate κ such that it is close enough to its optimal value (16),

$$\frac{1}{2}|q(m_\phi)\ell - N\kappa| \ll 1. \tag{18}$$

For an infinite number of plates this would allow to extend the mass range all the way to the frequency ω of the photons[8]. In practice, we can insert only a finite number of phase

[8]Above the photon frequency, ALP production is energetically forbidden and q becomes imaginary.

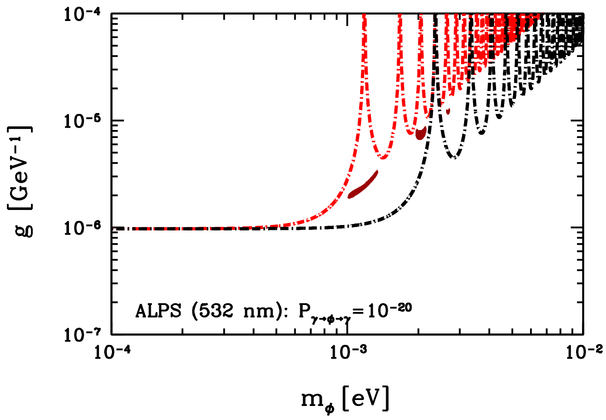

Figure 7: Iso-contours of the regeneration probability, as in Fig. 2. This figure demonstrates the potential for scanning through a whole range of masses by choosing the right κ for each m_ϕ. The red curve is the sensitivity without phase correction. The black curve is obtained by using three plates but scanning through a whole range of κ. In other words to obtain this curve one would insert the plates. Measure. Change the plates to a slightly different value of κ and measure again. This is repeated for all values of κ in the range $[0, 2\pi]$.

shift plates and Eq. (18) cannot be fulfilled for too large masses. But, already a small number of plates leads to a remarkable increase of the sensitivity for higher masses, as we can see from Fig. 7.

In summary: So called "light shining through a wall" experiments are a promising tool to search for light particles coupled to photons. In this note we have shown how the reach of such an experiment can be extended towards larger masses by inserting properly chosen phase shift plates. Although our explicit discussion is for the case of spin-0 axion-like particles the method works in general for particles exhibiting photon-particle-photon oscillations.

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