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A Cavity Experiment to Search for Hidden Sector Photons

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Abstract

We propose a cavity experiment to search for low mass extra U(1) gauge bosons with gauge-kinetic mixing with the ordinary photon, so-called paraphotons. The setup consists of two microwave cavities shielded from each other. In one cavity, paraphotons are produced via photon-paraphoton oscillations. The second, resonant, cavity is then driven by the paraphotons that permeate the shielding and reconvert into photons. This setup resembles the classic "light shining through a wall" setup. However, the high quality factors achievable for microwave cavities and the good sensitivity of microwave detectors allow for a projected sensitivity for photon-paraphoton mixing of the order of $\chi \sim 10^{-12}$ to 10^{-8}, for paraphotons with masses in the μeV to meV range – exceeding the current laboratory- and astrophysics-based limits by several orders of magnitude. Therefore, this experiment bears significant discovery potential for hidden sector physics.

Extensions of the standard model often contain extra U(1) gauge degrees of freedom. If the corresponding additional gauge bosons have direct renormalizable couplings to standard model matter, they are usually referred to as Z'-bosons. Negative collider searches for the latter have constrained their mass to $m_{Z'} \gtrsim$ few $\times\, 100\,\mathrm{GeV}$, for couplings of weak or electromagnetic strength [1].

However, in many cases, notably in realistic string-based scenarios standard model matter is uncharged under the additional U(1) symmetry and the corresponding gauge boson belongs to a "hidden sector", typically interacting with the standard model particles only via feeble gravity-like interactions. In these cases, the only renormalizable interaction with the standard model visible sector can occur via mixing [2; 3] of the photon γ with the hidden sector photon γ', often dubbed "paraphoton". Clearly, the sensitivity of collider experiments to photon mixing is extremely limited, in particular if the hidden sector photon has a small mass in the sub-eV range. Presently, the best laboratory limits on a low mass paraphoton and its mixing with the photon arise from Cavendish-type tests of the Coulomb law [4; 5] and from the search for signals of γ–γ' oscillations with laser "light shining through a wall" experiments [6; 7]. The best astrophysical limits come from considerations of the energy balance of stars, in particular the sun, and from the non-observation of photon regeneration in helioscopes [8; 9].

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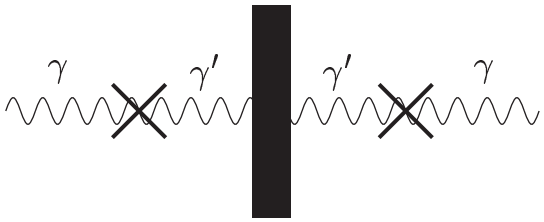

Figure 1: Schematic picture of a "light shining through a wall" experiment. The crosses denote the non-diagonal mass terms that convert photons into paraphotons. The photon γ oscillates into the paraphoton γ' and, after the wall, back into the photon γ which can then be detected.

In this letter, we propose a laboratory experiment to search for signatures of γ–γ' oscillations by exploiting high-quality microwave cavities. Our setup seems to be realizable with current technology and has a large window of opportunity for the discovery of low mass, $\mu\text{eV} \lesssim m_{\gamma'} \lesssim$ meV, hidden sector photons, exceeding the current limits on γ–γ' mixing by several orders of magnitude.

For definiteness, we will consider an extension of the standard model where one has, at low energies, say much below the electron mass, in addition to the familiar electromagnetic $\mathrm{U}(1)_{\mathrm{QED}}$, another hidden-sector $\mathrm{U}(1)_{\mathrm{h}}$ under which all standard model particles have zero charge. This may occur quite generally in string embeddings of the standard model (for general reviews, see e.g. Refs. [10; 11; 12; 13]), no matter whether they are based on compactifications of heterotic (e.g. [14; 15]), IIA (e.g. [16]), and IIB (e.g. [17]) string theory. The most general renormalizable Lagrangian describing these two U(1)'s at low energies is

$$\mathcal{L} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} - \frac{1}{4}B^{\mu\nu}B_{\mu\nu} - \frac{1}{2}\chi\, F^{\mu\nu}B_{\mu\nu} + \frac{1}{2}m_{\gamma'}^2 B_\mu B^\mu, \tag{1}$$

where $F_{\mu\nu}$ is the field strength tensor for the ordinary electromagnetic $\mathrm{U}(1)_{\mathrm{QED}}$ gauge field A^μ, and $B^{\mu\nu}$ is the field strength for the hidden-sector $\mathrm{U}(1)_{\mathrm{h}}$ field B^μ, i.e., the paraphoton. The first two terms are the standard kinetic terms for the photon and paraphoton fields, respectively. Because the field strength itself is gauge invariant for U(1) gauge fields, the third term is also allowed by gauge and Lorentz symmetry. This term corresponds to a non-diagonal kinetic term, a so-called kinetic mixing [3]. From the viewpoint of a low-energy effective Lagrangian, χ is a completely arbitrary parameter. Embedding the model into a more fundamental theory, it is plausible that $\chi = 0$ holds at a high-energy scale related to the fundamental theory. However, integrating out the heavy quantum fluctuations generally tends to generate non-vanishing χ at low scales. Indeed, kinetic mixing arises quite generally both in field theoretic [3] as well as in string theoretic [18; 19; 20; 21; 22; 23] setups. The last term in the Lagrangian (1) accounts for a possible mass of the paraphoton. This may arise from the breaking of the paraphoton $\mathrm{U}(1)_{\mathrm{h}}$ via a Higgs mechanism and choosing unitary gauge, or, alternatively, may be just an explicit Stückelberg mass term [24].

Let us now switch to a field basis in which the prediction of γ–γ' oscillations becomes apparent. In fact, the kinetic terms in the Lagrangian (1) can be diagonalized by a shift

$$B^\mu \to \tilde{B}^\mu - \chi A^\mu. \tag{2}$$

Apart from a multiplicative renormalization of the electromagnetic gauge coupling, $e^2 \to e^2/(1-\chi^2)$, the visible-sector fields remain unaffected by this shift and one obtains a non-diagonal mass term that mixes photons with paraphotons,

$$\mathcal{L} = -\frac{1}{4}F^{\mu\nu}F_{\mu\nu} - \frac{1}{4}\tilde{B}^{\mu\nu}\tilde{B}_{\mu\nu} + \frac{1}{2}m_{\gamma'}^2\left(\tilde{B}^\mu\tilde{B}_\mu - 2\chi\tilde{B}^\mu A_\mu + \chi^2 A^\mu A_\mu\right). \tag{3}$$

Therefore, in analogy to neutrino flavour oscillations, photons may oscillate in vacuum into paraphotons. These oscillations and the fact that the paraphotons do not interact with ordinary

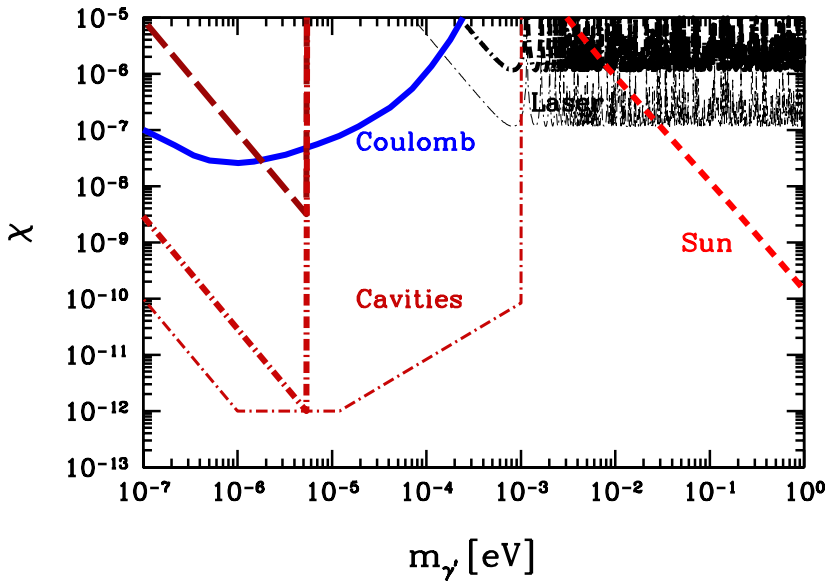

Figure 2: Existing bounds on the existence of massive paraphotons with kinetic mixing and projected sensitivity for the proposed experiment. Upper limit on the mixing parameter χ versus the mass $m_{\gamma'}$, obtained from the non-observation of deviations from Coulomb's law [4; 5] (blue, labelled "Coulomb"), from the non-observation of laser "light shining through a wall" (black, labeled "Laser"; thick: published limit from BFRT [6]; thin: projected sensitivity of ongoing experiments [7]), and from solar energy balance considerations [8; 9] (red, labelled "Sun"). Also shown is the projected sensitivity of our proposed "microwaves permeating through a shielding" setup (darkred, labelled "Cavities"). The dashed dotted line corresponds to the optimistic scenario ($Q = Q' = 10^{11}$, $\mathcal{P}_{\rm em} \sim 1$ W, $\mathcal{P}_{\rm detectable} = 10^{-26}$ W, $\nu_0 = 1.3\,$GHz, i.e. $\omega_0 \approx 5.4\,\mu$eV) and the dashed fat line to the more modest one ($Q = 10^{10}$, $Q' = 10^4$, $\mathcal{P}_{\rm em} \sim 1$ W, $\mathcal{P}_{\rm detectable} = 10^{-20}$ W, $\nu_0 = 1.3$ GHz) in the text. In both cases we have used $|G| = 1$ for $m_{\gamma'} \leq \omega_0$ and $|G| = 0$ for $m_{\gamma'} > \omega_0$, for simplicity, for the "geometry factor" (25). The thin dashed dotted line corresponds to the sensitivity which one might get from the optimistic scenario, if one scans the frequency between 250 MHz $\lesssim \nu_0 \lesssim$ 250 GHz, corresponding to $1\,\mu\text{eV} \lesssim \omega_0 \lesssim 1\,$meV (for frequencies $\nu_0 > 3\,$GHz, the losses in the cavities grow due to an increased surface resistance [26]; accordingly, we have assumed a drop in the Q value for frequencies higher than 3 GHz.)

matter forms the basis of the possibility [2] to search for signals of paraphotons in "light shining through a wall" experiments (cf. Fig. 1). The sensitivity of ongoing experiments of this type (for a review, see Ref. [25]) for paraphoton searches has recently been estimated in Ref. [7]. Their discovery potential extends the current upper limit on χ set by the BFRT collaboration [6] by about one order of magnitude over the whole range of masses $m_{\gamma'}$ (cf. Fig. 2).

Here, we propose another setup searching for signatures of γ–γ' oscillations which resembles the classic "light shining through a wall" setup. It consists of two microwave cavities shielded from each other (cf. Fig. 3). In one cavity, paraphotons are produced via photon-paraphoton oscillations. The second, resonant, cavity is then driven by the paraphotons that permeate the shielding and reconvert into photons. Due to the high quality factors achievable for microwave cavities and the good sensitivity of microwave detectors such a setup will allow for an unprecedented discovery potential for hidden sector photons in the mass range from μeV to meV range (cf. Fig. 2).

Before we start with a detailed calculation, let us present a simple estimate based on a comparison with the familiar "light shining through a wall" setup, which exploits an optical

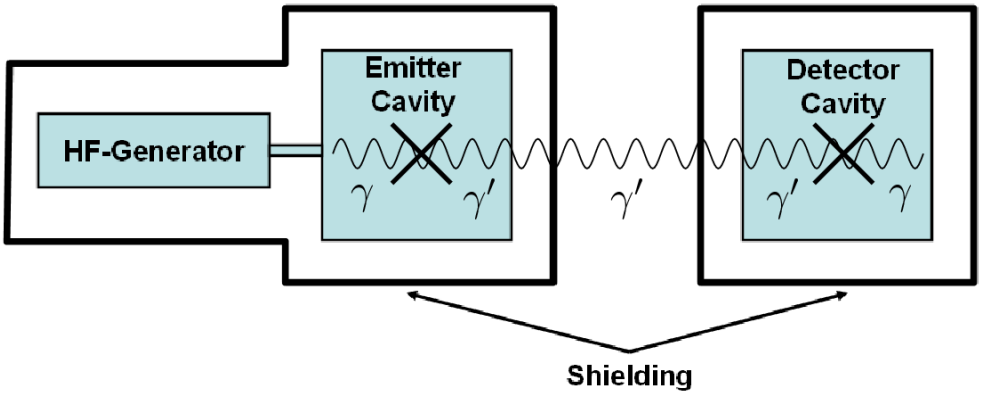

Figure 3: Schematic illustration of a "microwaves permeating through a shielding" experiment for the search for massive hidden sector photons mixing with the photon (a high-frequency (HF) generator drives the emitter cavity).

cavity and a laser with wavelength $\sim$ few $\times$ 100 nm. In optical cavities, the spatial extent of the laser beam transverse to the beam direction ($\sim$ mm – cm) is much greater than the wavelength. The wave is effectively a plane wave propagating in the beam direction and the problem is effectively one-dimensional. This is not the case for microwave or radio-frequency (RF) cavities where the size of the cavity is similar to the wavelength in all three directions. Nevertheless, let us for the moment imagine an unrealistic cavity which has infinite extent in two directions. Then the situation is equivalent to a standard "light shining through a wall" experiment (the shielding is equivalent to a wall). For a setup with cavities on both sides of the wall, the probability for a photon to pass through the wall and to be emitted by the second cavity is [2; 7; 27; 28],

$$P_{\rm trans} = 16\chi^4 \left[\frac{N_1+1}{2}\right] \left[\frac{N_2+1}{2}\right] \left[\sin\left(\Delta k \ell_1\right) \sin\left(\Delta k \ell_2\right)\right]^2 . \tag{4}$$

Here, $N_{1,2}$ are the number of passes the light makes through the cavities, $\ell_{1,2}$ are the lengths of the two cavities, and

$$\Delta k = \omega - \sqrt{\omega^2 - m_{\gamma'}^2} \tag{5}$$

is the momentum difference between the photon and the paraphoton, expressed in terms of the energy of the laser photons, $\omega = 2\pi\nu$, where ν is the frequency of the laser light. Maximal sensitivity to the mixing parameter χ will be achieved if both sines in Eq. (4) are equal to one. One way to achieve this is to choose the angular frequency $\omega \equiv 2\pi/\lambda = m_{\gamma'}$ and in exploiting cavities of length $\ell_1 = \ell_2 = \lambda/2$, where λ is the wavelength of the laser light.

Using this we can get a rough idea of what may be accomplished by a similar experiment using microwave or RF cavities (cf. Fig. 3) instead of optical cavities. Using $(N+1)/2 \sim Q$, where Q is the quality factor of the cavity, we roughly expect

$$P_{\rm trans}^{\rm max} \sim \chi^4 Q Q'. \tag{6}$$

To get an idea of the sensitivity which such an experiment can reach let us plug in some numbers. The power output $\mathcal{P}_{\rm out}$ of the detector cavity[1] will be

$$\mathcal{P}_{\rm out} = P_{\rm trans} \mathcal{P}_{\rm in}, \tag{7}$$

[1]When we speak of power going into and out of the cavities we can alternatively think of this as a measure for the energy stored inside the cavity. The power is related to the stored energy U and the Q factor via $\mathcal{P} = \omega U/Q$.

in terms of the power $\mathcal{P}_{\rm in}$ put into the emitter cavity. An input power of $\mathcal{P}_{\rm in} \sim 1$ W is quite realistic[2] [30] and an emission of $\mathcal{P}_{\rm out} \sim 10^{-26}$ W is just on the verge of being detectable [31]. High quality cavities based on superconducting technology can reach $Q \sim 10^{11}$ [29]. Plugging these numbers into Eqs. (6) and (7), we infer that, very optimistically, a setup based on microwave or RF cavities might be sensitive to values of the mixing parameter as small as $\chi \sim 10^{-12}$, in the mass range corresponding to the frequency range, i.e. $\mu\text{eV} \lesssim m_{\gamma'} \lesssim \text{meV}$.

Motivated by this estimate, let us proceed to more realistic situations, taking into account the appropriate fully three-dimensional geometry. Our starting point are the equations of motion for the photon field A and the paraphoton field B (Lorentz indices suppressed[3]) following from Eq. (1),

$$(\partial^\mu \partial_\mu + \chi^2 m_{\gamma'}^2) A = \chi m_{\gamma'}^2 B, \tag{8}$$

$$(\partial^\mu \partial_\mu + m_{\gamma'}^2) B = \chi m_{\gamma'}^2 A. \tag{9}$$

Our strategy is as follows (cf. Fig. 3). We start with the ordinary electromagnetic field inside the first "emitter" cavity. This field acts as a source for the paraphoton field. The paraphoton field permeates the shielding and in turn acts as a source for an electromagnetic field inside the second "detector" cavity. We will always consider the lowest non-trivial order.

To lowest order in χ, we can obtain the electromagnetic field inside the emitter cavity by solving $\partial^\mu \partial_\mu A = 0$, i.e. the standard equation of ordinary electrodynamics. Implementing the (time independent) boundary conditions of the cavity is a textbook exercise [32]. Using the separation ansatz

$$A_{\rm em}(\mathbf{x}, t) = a_{\rm em}(t) A_{\omega_0}(\mathbf{x}), \tag{10}$$

accounting for a finite quality factor Q of the cavity, and including a driving force $f(t)$, we have,

$$\left(\frac{d^2}{dt^2} + \frac{\omega_0}{Q} \frac{d}{dt} + \omega_0^2 \right) a_{\rm em}(t) = f(t), \tag{11}$$

where

$$-\nabla^2 A_{\omega_0}(\mathbf{x}) = \omega_0^2 A_{\omega_0}(\mathbf{x}) \tag{12}$$

is an eigenfunction of the spatial part of the wave equation including the appropriate boundary conditions. It is convenient to choose a normalization,

$$\int_V d^3\,\mathbf{x} |A_{\omega_0}(\mathbf{x})|^2 = 1. \tag{13}$$

For example, if the cavity is a cube with side length L, the eigenfunctions for the electric field in the z-directions are

$$A_{\omega_0^{mnp}}(\mathbf{x}) = C_{mnp} \sin\left(\frac{m\pi x}{L}\right) \sin\left(\frac{n\pi y}{L}\right) \cos\left(\frac{p\pi z}{L}\right), \tag{14}$$

where C_{mnp} are normalization factors. The eigenvalues are in this case given by

$$\omega_0^{mnp} = \frac{\pi}{L}\sqrt{m^2 + n^2 + p^2}, \qquad m, n = 1, 2, \ldots \qquad p = 0, 1, \ldots . \tag{15}$$

Employing an oscillating driving force,

$$f(t) = f_0 \exp(-\mathrm{i}\omega t), \tag{16}$$

[2]Let us express this in terms of the energy stored inside the cavity. For example at the frequency 1.3 GHz used in the TESLA cavities and with a $Q \sim 10^{11}$ [29] this corresponds to an energy $U \sim 10$ J stored inside the cavity.

[3]Although we may think of A as the gauge potential. Using Coulomb gauge $A^0 = 0$, which is compatible with Lorentz gauge, we can immediately relate A to electric fields via $E = -\frac{dA}{dt}$.

the amplitude will eventually approach,

$$a_{\rm em}(t) = a^0_{\rm em}\exp(-{\rm i}\omega t) = \frac{f_0}{\omega^2-\omega_0^2-{\rm i}\frac{\omega\omega_0}{Q}}\exp(-{\rm i}\omega t). \tag{17}$$

For $Q \gg 1$ and a driving force that is resonant with the cavity, $\omega = \omega_0$, the amplitude is enhanced by a factor of Q with respect to the driving force,

$$a^0_{\rm em} \approx {\rm i}\frac{Q}{\omega_0^2}f_0. \tag{18}$$

The field $a_{\rm em}(t)A_{\omega_0}(\mathbf{x})$ now acts as a source on the right hand side of the equation of motion for the paraphoton fields, Eq. (9). The paraphoton does not interact with ordinary matter and no boundary conditions are enforced at finite $\mathbf{x}$. The appropriate solution are therefore obtained from the (retarded) massive Greens function,

$$B(\mathbf{x},t) = \chi m^2_{\gamma'}\int_V d^3\mathbf{y}\frac{\exp({\rm i}k|\mathbf{x}-\mathbf{y}|)}{4\pi|\mathbf{x}-\mathbf{y}|}a_{\rm em}(t)A_{\omega_0}(\mathbf{y}), \tag{19}$$

where V is the volume of the emitter cavity and

$$k^2 = \omega^2 - m^2_{\gamma'}. \tag{20}$$

In our detector cavity, the field $B(\mathbf{x},t)$ now acts as a source, i.e. a driving force. The wave equation can again be solved by a separation ansatz analog to Eq. (10), $A'(\mathbf{x},t) = a_{\rm det}(t)A'_{\omega'_0}(\mathbf{x})$,

$$\left(\frac{d^2}{dt^2} + \frac{\omega'_0}{Q'}\frac{d}{dt} + \omega_0'^2\right)a_{\rm det}(t) = b(t). \tag{21}$$

The driving force $b(t)$ can be obtained by remembering that the spatial eigenfunctions of cavities form a complete orthonormal set. Inserting the separation ansatz into Eq. (9), multiplying by the eigenfunction $A'_{\omega'_0}(\mathbf{x})$ and integrating over the volume V' of the detector cavity, we find

$$b(t) = \chi^2 m^4_{\gamma'}a_{\rm em}(t)\int_{V'}\int_V d^3\mathbf{x}\,d^3\mathbf{y}\,\frac{\exp({\rm i}k|\mathbf{x}-\mathbf{y}|)}{4\pi|\mathbf{x}-\mathbf{y}|}A_{\omega_0}(\mathbf{y})A'_{\omega'_0}(\mathbf{x}). \tag{22}$$

To get resonant enhancement we choose

$$\omega'_0 = \omega_0. \tag{23}$$

The integral in Eq. (22) has dimensions length^2 = frequency^{-2}. Taking this into account we write

$$b(t) = a_{\rm em}(t)\frac{\chi^2 m^4_{\gamma'}}{\omega_0^2}G(k/\omega_0) \tag{24}$$

where G is a dimensionless function that encodes the geometric details of the setup, e.g. relative position, distance and shapes of the cavities. Moreover, it depends on the mass $m_{\gamma'}$ via k,

$$G(k/\omega_0) \equiv \omega_0^2\int_{V'}\int_V d^3\mathbf{x}\,d^3\mathbf{y}\,\frac{\exp({\rm i}k|\mathbf{x}-\mathbf{y}|)}{4\pi|\mathbf{x}-\mathbf{y}|}A_{\omega_0}(\mathbf{y})A'_{\omega_0}(\mathbf{x}). \tag{25}$$

It is typically of order one, as can be seen from Fig. 4, where we show G for a setup with two identical cubic cavities.

After some time the amplitude in the detector cavity will approach

$$a^0_{\rm det} = {\rm i}Q'\frac{\chi^2 m^4_{\gamma'}}{\omega_0^4}G\,a^0_{\rm em}. \tag{26}$$

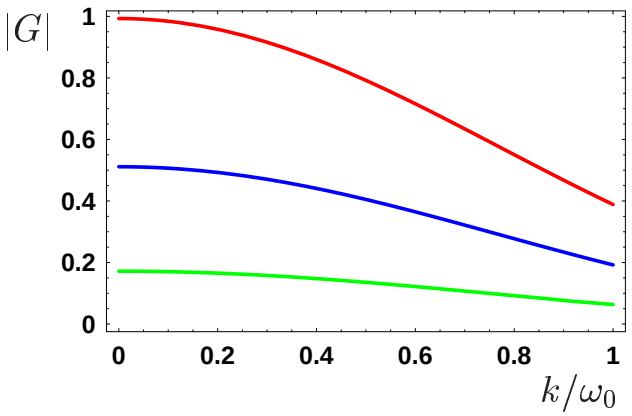

Figure 4: Geometry factor $|G|$ for a setup with two identical cubic cavities with side length $L = \sqrt{2}\pi/\omega_0$ in the $n = 1, m = 1, p = 0$ mode of Eq. (14). The cavities are placed parallel and are separated by a distance $d = 0$ (red), $d = L$ (blue) and $d = 5\,L$ (green) along the z-axis. As expected $|G|$ scales roughly with $1/d$.

Finally, we have to relate the amplitudes to the power input/output in the emitter/detector cavities. The quality factor is directly related to the power consumption/emission of a cavity,

$$\mathcal{P} = \frac{\omega_0}{Q} U, \tag{27}$$

where

$$U = const\,|a|^2 \tag{28}$$

is the energy stored inside the cavity. Using this relation, the probability for a photon to pass through the shielding and to be emitted by the second cavity is

$$P_{\rm trans} = \frac{\mathcal{P}_{\rm det}}{\mathcal{P}_{\rm em}} = \frac{Q}{Q'} \frac{|a^0_{\rm det}|^2}{|a^0_{\rm em}|^2} = \chi^4\, QQ'\, \frac{m^8_{\gamma'}}{\omega^8_0}\, |G|^2. \tag{29}$$

If we choose the cavity frequency to be $\omega_0 = m_{\gamma'}$, our expression agrees up to a factor of order unity with our estimate (6).

Let us now turn to the $m_{\gamma'}$ dependence of the effect. As can be seen from Fig. 4, the k and therefore $m_{\gamma'}$ dependence of $|G|$ is not very strong. Moreover, the latter is non-zero[4] for all allowed $\omega_0 \geq k \geq 0$. Therefore, the sensitivity of the proposed experiment to the value of the mixing parameter χ decreases roughly as $\omega_0^2/m_{\gamma'}^2$ when going to smaller masses. What happens for $m_{\gamma'} > \omega_0$? In this case $k = {\rm i}\kappa$ is imaginary and $|G|$ drops exponentially as $|G| \sim \exp(-\kappa d)$ where d is the distance between the cavities. Since some distance between the cavities is necessary to allow for shielding etc., the experiment will not be very sensitive in this region.

In Fig. 2, we sketch the sensitivity region of two scenarios:

- An optimistic scenario where we basically stick together the best cavities $Q = Q' \sim 10^{11}$, $\mathcal{P}_{\rm em} \sim 1$ W, best detectors $\mathcal{P}_{\rm detectable} \sim 10^{-26}$ W, and assume that perfect shielding is possible and
- a more modest scenario where we use $Q \sim 10^{10}$, $Q' \sim 10^4$, $\mathcal{P}_{\rm em} \sim 1$ W, and a detectable power of $\mathcal{P}_{\rm detectable} \sim 10^{-20}$ W.

[4]It should be mentioned that this is typical for the lowest cavity modes where the field does not change sign inside the cavity.

In both cases the achieved sensitivity is better than the current laboratory *and* astrophysical limits. The sensitivity region can be further extended by performing several experiments at different frequencies or, even better, scanning through a whole region of frequencies (thin dashed dotted line in Fig. 2).

At last we note that one can also obtain a bound from the observed maximal Q value of the emitter cavity itself. The conversion of photons into paraphotons leads to an energy loss in the emitter cavity. From our discussion above we can estimate that the probability for a photon to convert to a paraphoton during one pass through the cavity is

$$P_{\rm loss} \sim \chi^2 \frac{m_{\gamma'}^4}{\omega_0^4}. \tag{30}$$

Assuming that conversion into paraphotons is the only source of energy loss we infer that the maximal possible Q is

$$Q^{\rm max} \sim \frac{2\pi}{P_{\rm loss}} \sim \frac{2\pi}{\chi^2}\frac{\omega_0^4}{m_{\gamma'}^4}. \tag{31}$$

At the resonance frequency, $\omega_0 = m_{\gamma'}$, an observed value $Q > 10^{11}$ will then translate into a bound of roughly $\chi(m_{\gamma'} = \omega_0) \lesssim 10^{-5}$. As above, the bound becomes weaker as $\sim \omega_0^2/m_{\gamma'}^2$, for smaller $m_{\gamma'}$, and drops off sharply for $m_{\gamma'} > \omega_0$.

Finally, let us comment on a few experimental issues. First, since we use cavities both for production and detection we have to assure that both cavities have the same resonant frequency. More precisely the frequencies have to agree in a small range $\Delta\omega_0/\omega_0 \sim 1/Q'$. This is a non-trivial task. However, compared to optical frequencies (as proposed in [27; 28]), this should be significantly simpler for microwave of RF cavities: the wavelength is longer and correspondingly the allowed inaccuracies in the cavity are much larger. Indeed, the cavities originally developed for the TESLA accelerator [29] may be mutually tuned in frequencies to a few $\times$ 100 Hz [30]. With a resonance frequency of roughly 1 GHz, this corresponds to an allowed quality factor of the detector cavity of $Q' \sim 10^6$. We have used even a somewhat smaller $Q' \sim 10^4$ in our modest scenario. Second, one needs to provide sufficiently good shielding between the cavities to prevent exciting the detector cavity by ordinary electromagnetic fields leaking from the production region. This is closely linked to the question how one can decide that a possible signal is physical in origin. One way to accomplish this could consist in checking the phase of a "signal". The phase differs between an artifact resulting from a ordinary photon sneaking out of the cavity and a true paraphoton signal: for a true signal the phase is encoded in the complex phase of G. The photon is massless and the wavenumber is $k_\gamma = \omega_0$ whereas the paraphoton is massive and has a smaller wavenumber $k = \sqrt{\omega_0^2 - m_{\gamma'}^2} < \omega_0 = k_\gamma$. Therefore, the phase difference, $\Delta\phi$, between an artifact and a true signal is approximately

$$\Delta\phi \sim (k_\gamma - k)d = (\omega_0 - k)d, \tag{32}$$

where d is the distance between the cavities[5].

Last, but not least, let us note that the experimental setup proposed in this letter can be extended [34] to a search facility for light neutral spin-zero (axion-like) particles ϕ, coupling to electromagnetism, at low energies, according to

$$\mathcal{L} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} + \frac{1}{2}\partial_\mu\phi\partial^\mu\phi - \frac{1}{2}m_\phi^2\phi^2 - \frac{1}{4}g\phi F_{\mu\nu}\tilde{F}^{\mu\nu}\,, \tag{33}$$

[5]One might also wonder how one can distinguish this signal from a signal caused by an electric current of minicharged particles, the latter being Schwinger pair produced in the electric field of the cavity, as suggested in Ref. [33]. This is actually quite simple. Since such a current would flow in the direction of the electric field one can simply choose the separation between the two cavities in our setup to be perpendicular to the electric field. The minicharged particle current would then simply miss the detector cavity.

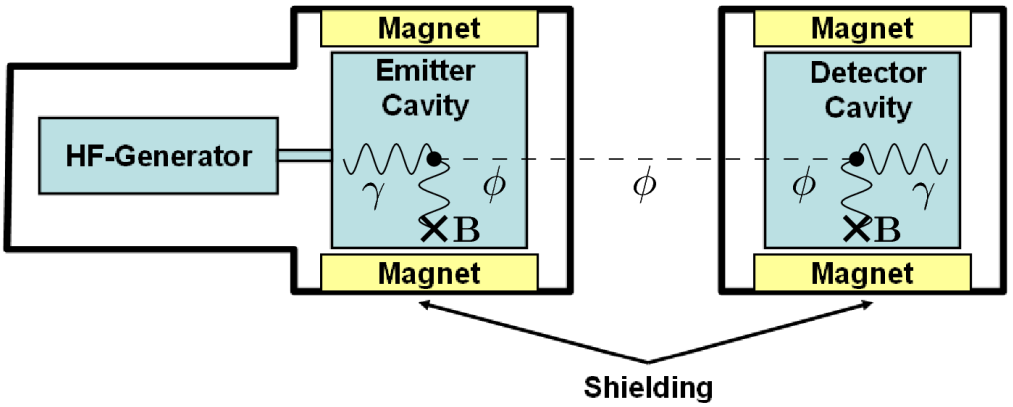

Figure 5: Schematic illustration of a "microwaves permeating through a shielding" experiment for the search for an axion-like particle ϕ mixing with the photon in the presence of a magnetic field.

where $\tilde{F}_{\mu\nu}$ is the dual electromagnetic field strength tensor[6]. In fact, by placing both the emitter as well as the detector cavity each into a magnet of strength B, one may drive the detector cavity now with the axion-like particles which have been produced in the emitter cavity and which have reconverted in the detector cavity (cf. Fig. 5). From a calculation similar to the one presented in this letter, one finds, in analogy to (29), for the probability that a photon passes through the shielding and is emitted by the second cavity [34],

$$P_{\rm trans} \sim \left(\frac{g\,B}{\omega_0}\right)^4 QQ'\,|G|^2. \tag{34}$$

Let us estimate the discovery potential of such an experiment, given current technology. State-of-the-art axion dark matter experiments such as ADMX [31] exploit standalone RF cavities based on normal-conducting technology[7] with $Q \sim 10^6$ inside a strong magnet with $B \sim 10$ T. Using these numbers and $P_{\rm em} = 1$ W, $P_{\rm det} = 10^{-26}$ W, we estimate a sensitivity of $g \sim 8 \times 10^{-10}$ GeV^{-1}, for $m_\phi \lesssim \omega_0 = 5$ μeV – about one order of magnitude above the (albeit model-dependent, cf., e.g., [35; 36]) limits set by solar energy loss considerations [37] and by the non-observation of solar axion-like particle induced photon regeneration by the CAST collaboration [38], but considerably better than the limits of present day optical "light shining through a wall" experiments (for a review and references, see, e.g., Ref. [25]). An improvement in the sensitivity to the range $g \sim (10^{-15}$ to $10^{-14})$ GeV^{-1}, predicted for proper QCD axions [39; 40; 41] in the μeV mass range [42; 43; 44], will require still substantial technical advances.

In conclusion: In this letter, we have proposed a simple experiment to search for massive hidden sector photons that have kinetic mixing with ordinary photons. The experiment would allow to probe a region of parameter space that is so far unexplored by laboratory experiments *as well as* astrophysical observations. Therefore, it bears significant discovery potential for hidden sector physics.

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[6]The effective Lagrangian (33) applies for a pseudoscalar ϕ, corresponding to a parity odd spin-zero boson. In the case of a scalar axion-like particle, the $F_{\mu\nu}\tilde{F}^{\mu\nu}$ in Eq. (33) is replaced by $F_{\mu\nu}F^{\mu\nu}$.

[7]In the strong magnetic field required for this setup it is unclear whether one can use a superconducting cavity that would, in principle, allow for a higher Q value.

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