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Review

Perspective to search for sub-eV neutral boson resonances with stimulated laser colliders

K. Homma[1,2,a]

[1] Graduate School of Science, Hiroshima University, Kagamiyama, Higashi-Hiroshima, Hiroshima 739-8526, Japan
[2] IZEST, Ecole Polytechnique, 91128 Palaiseau, France

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Abstract. We present a novel approach to search for sub-eV neutral bosons coupling to two photons which can be candidates for dark fields in the universe. Similarly to conventional particle colliders, the search aims at detection of resonance states of these dark fields produced by photon-photon scattering, however, the interaction rate is enhanced by the stimulated decay of resonances into two photons under a background coherent laser field. We discuss the future experimental strategy to explore a large mass-coupling parameter space with high-intensity lasers.

1 Introduction

Understanding a large fraction of dark components in the universe is one of the most fundamental problems in moder physics. In high energy scales, several neutral bosons have been discovered. It is well known that the neutral pion and the Higgs-like particle can couple to two photons via the decay processes at 100 MeV and 100 GeV energy scales, respectively. These facts encourage further experimental searches for similar type of fields via two-photon coupling in very different energy scales in general. For instance, there are theoretical rationales to expect sub-eV particles such as the axion [1–3] (pseudoscalar boson) and the dilaton [4] (scalar boson) associated with breaking of fundamental symmetries in the context of particle physics and cosmology. Furthermore, the advent of high-intensity laser systems and the rapid leap of the intensity encourage the approach to probe weakly coupling dark fields with optical photons by the enhanced luminosity factor. With high-intensity lasers, we might be able to unveil the different aspects of the quantum vacuum at different space-time scales based on analogous observables in quantum optics. We present the novel approach to realize the laboratory search for low-mass and weakly coupling dark fields which can be dark components of the universe by detecting four-wave mixing of two-color laser fields in the vacuum. This can be interpreted as a kind of quasi-parallel photon-photon collider whose interaction rate is enhanced by the resonant production of a sub-eV neutral boson and also by the stimulated decay in the coherent laser fields [5,6]. We review these two important aspects of the novel approach. We

[a] e-mail: homma@hepl.hiroshima-u.ac.jp

then discuss the experimental strategy to search for sub-eV dark fields with high-repetition rate and high-intensity laser systems in the large mass-coupling parameter space.

2 Dynamics of photon-photon scattering

We discuss photon-photon scattering with the effective Lagrangian in Eq. (1) [7] as generic as possible,

$$-L_\phi = gM^{-1}\frac{1}{4}F_{\mu\nu}F^{\mu\nu}\phi, \quad -L_\sigma = gM^{-1}\frac{1}{4}F_{\mu\nu}\tilde{F}^{\mu\nu}\sigma, \tag{1}$$

where an effective coupling between two photons and a dark field scalar ϕ or pseudoscalar σ is represented by g^2/M. For instance, if the dimensional constant M is vacuum expectation value 246 GeV and the dimensionless coupling g^2 is $\sim \alpha_{qed}$, the effective coupling corresponds to that of Higgs-like particle to two photons. If we are based on the invisible axion scenario [5,6], a dark field satisfying the dimensional constant $M = 10^{11} - 10^{16}$ GeV and the mass m = meV $-$ μeV can be cold dark matter candidates. If M corresponds to the Planck mass $M_P \sim 10^{18}$ GeV, on the other hand, the interaction is as weak as that of gravity and this case would have a great relevance to Dark Energy if $m \geq$ neV [4].

We may discuss possibilities to exchange scalar and pseudoscalar type of fields by requiring combinations of photon polarizations in the initial and final states [8]. The virtue of laser experiments is in the specifications on the all photon spin states both in the initial and final states in the two body photon-photon interaction. This allows us to distinguish types of exchanged fields in general.

3 Inclusion of a resonance state in Quasi-Parallel photon-photon colliding System (QPS)

Because we aim at extremely low-mass ranges below eV, we have to reduce the center-of-mass system (CMS) energy of photon-photon collisions. For this purpose, we introduce quasi-parallel colliding system (QPS) with the optical laser energy of $\sim$1 eV. The CMS energy with variables in QPS can be defined as

$$E_{cms} = 2\sin\vartheta\omega, \tag{2}$$

where ϑ is defined as a half incident angle between two incoming photons with $\vartheta \ll 1$ and ω is the beam energy in unit of $\hbar = c = 1$. This relation indicates that we have two experimental knobs to adjust E_{cms}. If we take the head-on collision geometry, we have to introduce very long wavelength as the incident photons. However, it is comparatively easier to introduce a small incident angle by focusing a single laser beam. In such a case E_{cms} can be lowered by keeping ω constant. The feature of QPS is that decayed photons experience frequency shifts and detectable signature of the interaction are produced, because QPS corresponds to the system boosting CMS into the direction perpendicular to the colliding axis. The frequency shift is prominent along that boost axis. In the forward direction on the boost axis we expect a frequency up shift to close to the double of the incident frequency, while a zero frequency photon must be emitted to the backward direction due to the energy-momentum conservation independent of the dynamics of the exchanged field. This may be interpreted as if the second harmonic photon is generated from the nonlinear vacuum response. This could be an interesting analogue to second harmonic generation due to the nonlinear response of a crystal with a laser injection which was pioneered by Franken et al. [9].

We focus on the direct production of a resonance state via s-channel Fayman amplitude in photon-photon collisions in QPS, because if the exact resonance condition $m = E_{cms}$ is satisfied, the peak amplitude becomes independent of the weakness of the coupling as shown below. The square of the scattering amplitude A can be expressed as Breit-Wigner (BW) resonance function [10]

$$|A|^2 = (4\pi)^2 \frac{W^2}{\chi^2(\vartheta) + W^2}, \tag{3}$$

where χ and the width W are defined as $\chi(\vartheta) \equiv \omega^2 - \omega_r^2(\vartheta)$ and $W \equiv (\omega_r^2/16\pi)(g^2 m/M)^2$, respectively. The energy ω_r satisfying the resonance condition can be defined as $\omega_r \equiv m^2/(1-\cos 2\vartheta_r)$ [11]. If this condition is satisfied, $\chi^2(\vartheta_r) = 0$ and thus $|A|^2$ approaches to $(4\pi)^2$. If we take M as the Planckian mass, however, the width W becomes extremely small. This implies that the resonance width is too small to hit the peak of the resonance function. How can we overcome this situation? The special nature of QPS improves this situation. In the focused laser field, the incident angle becomes uncertain due to the wavy nature of photons. This implies that $|A|^2$ must be averaged over the possible uncertainty on E_{cms}. This unavoidable integration over the possible angular uncertainty provides $1/M^2$ dependence of $|A|^2$ compared to the $1/M^4$ dependence when no resonance is contained in the uncertainty range, that is, when $\chi^2(\vartheta) \gg W^2$.

We expect that the cross section of photo-photon scattering in QED process σ_{qed} in QPS is quite suppressed due to the sixth power dependence on the CMS energy and the fourth power dependence on the incident angle which is expressed as $\sigma_{qed} \sim (\alpha^2/m_e^4)^2 \omega^6 \vartheta^4$ with the fine structure constant α and electron mass m_e [12]. Therefore, the low frequency photons in QPS is one of the best systems to probe sub-eV dark fields.

4 Stimulated decay by coherent field and four-wave mixing

The inclusion of a resonance state within the uncertainty on E_{cms} is still short in order for us to reach the sensitivity to gravitational coupling strength. We thus considered the stimulation of Faynman amplitude by replacing the quantum vacuum state $|0\rangle$ with the quantum coherent state $|N\gg$ as follows [7]. A laser field is represented by the quantum coherent state which features a superposition of different photon numbers, characterized by an averaged photon number N [13],

$$|N\gg \equiv \exp(-N/2) \sum_{n=0}^{\infty} \frac{N^{n/2}}{\sqrt{n!}} |n>, \tag{4}$$

where $|n>$ is the normalized state of n photons

$$|n> = \frac{1}{\sqrt{n!}} \left(a^\dagger\right)^n |0>, \tag{5}$$

with $a^\dagger$ and a the creation and the annihilation operators of photons specified with momentum and polarization, respectively. The coherent state satisfies the normalization condition

$$\ll N|N\gg = 1. \tag{6}$$

We also derive following properties of coherent states $|N\gg$ and $\ll N|$:

$$a|N\gg = \sqrt{N}|N\gg \text{ and } \ll N|a^\dagger = \sqrt{N}\ll N| \tag{7}$$

from the familiar relations

$$a^\dagger|n\rangle = \sqrt{n+1}|n+1\rangle \text{ and } a|n+1\rangle = \sqrt{n+1}|n\rangle. \tag{8}$$

The property in Eq. (7) gives the expectation value of the annihilation and creation operators to coherent states

$$\ll N|a|N\gg = \sqrt{N} \text{ and } \ll N|a^\dagger|N\gg = \sqrt{N}. \tag{9}$$

We now see how the coherent states are related with the Feynman amplitude for the process $p_1 + p_2 \to p_3 + p_4$ by the exchange of a dark field with the mass m as follows

$$V_2\left((p_1+p_2)^2 + m\right)^{-1} V_1, \tag{10}$$

where V_1 and V_2 correspond to the matrix elements of the interaction defined in Eq. (1).

$$\begin{aligned} V_1 &= gM^{-1}{<}0|F_{\mu\nu}|1_{p_1}{>}{<}0|F^{\mu\nu}|1_{p_2}{>} \quad \text{and} \\ V_2 &= gM^{-1}{<}1_{p_3}|F_{\mu\nu}|0{>}{<}1_{p_4}|F^{\mu\nu}|0{>}, \end{aligned} \tag{11}$$

where $|1_{p_i}{>}$ with $i = 1, 2$ denotes one-photon states specified with the individual momentum p_i. In order to grasp the essence, we abbreviate the factors included in the initial vertex as follows by suppressing all the complications arising from the momenta and polarization vectors

$${<}0|F_{\mu\nu}|1_{p_i}{>} \to {<}0|a|1_{p_i}{>} = 1, \tag{12}$$

where the second of Eq. (8) is applied to $n = 0$. We now consider the case where p_i is annihilated within the coherent sea with the same momentum and polarization state, and the proceeding interaction still occurs in the presence of the coherent field. We then extend the initial state from the one-photon state $|1_{p_i}\rangle$ to the coherent states $|N_{p_i}\gg$

$$\ll N_{p_i}|a|N_{p_i}\gg = \sqrt{N_{p_i}}, \tag{13}$$

where the vacuum state $\langle 0|$ is replaced by $\ll N_{p_i}|$. This corresponds to the result of the first of Eq. (9). We can see that the presence of the coherent states enhances the matrix element V_1 by $\sqrt{N_{p_1}}\sqrt{N_{p_2}}$. As for the second vertex V_2, the situation is different. The outgoing photons with the momenta p_3 and p_4 are not included in the initial ones, because we require p_3 to be the observational signal with frequency up-shift. Thus, the final state photons are created spontaneously from the vacuum state $|0\rangle$ with no enhancement, because the sea of the coherent state specified with p_3 and p_4 does not exist. If $p_1 = p_2$ is realized by simplifying that the incident two photons are annihilated in a single-frequency mode laser with the averaged number of photons N, the interaction rate is proportional to the square of the Feynman amplitude. Therefore, we obtain the enhancement factor

$$(\sqrt{N}\sqrt{N})^2 = N^2. \tag{14}$$

If we could supply $|N_{p_4}\gg$ in advance of the interaction, the situation changes. This additional laser beam, referred to as inducing laser, provides us with a sea of photons from which the photon p_4 is created in the induced manner, so that $<p_4|F^{\mu\nu}|0>$ in the second of Eq. (11) will be modified to

$$<1_{p_4}|a^\dagger|0> = 1 \to \ll N_{p_4}|a^\dagger|N_{p_4}\gg = \sqrt{N_{p_4}}. \tag{15}$$

This is the time-reversed process of the ones for the incident photons with p_1 and p_2 in V_1. We note that photon p_3 remains to be created spontaneously from the vacuum because we must be able to detect it as the signature of the interaction, while no attempt is made to observe the photon p_4, which will be embedded quietly in $\ll N_{p_4}|$. We finally reach an important consequence. The overall enhancement factor on the interaction rate is expressed as

$$(\sqrt{N_{p_1}}\sqrt{N_{p_2}}\sqrt{N_{p_4}})^2 = N_{p_1}N_{p_2}N_{p_3}. \tag{16}$$

This enhancement factor is similar to the known atomic process in quantum optics. If we assign photon energies as $|p_1| = |p_2| = \omega$, $|p_4| = u\omega$ and $|p_3| = (2-u)\omega$ with $0 < u < 1$, the following transition is expected in the vacuum via energy-momentum conservation

$$\omega + \omega \to \phi/\sigma \to (2-u)\omega + u\omega. \tag{17}$$

This is the similar transition discussed in Coherent Anti-Stokes Raman Spectroscopy (CARS) [14], more generally, four-wave mixing, where the role of the resonance production and decay in the vacuum is replaced by the nonlinear atomic process.

Because N_{p_i} has no limitation due to the bosonic nature of photons, we can expect a huge enhancement factor by the cubic dependence on the photon numbers. In conventional charged particle colliders, on the other hand, the dependence on the number of particles is quadratic and also there is a physical limitation by the space charge effect. Therefore, the stimulated photon collider would provide an extraordinary sensitivity to weak couplings compared to that of particle colliders whose prime missions is, of course, to produce heavy new particles.

5 Experimental strategy

Given the concept of stimulated photon collisions in QPS, the key control parameters or experimentally adjustable knobs are summarized below.

The resonance condition is satisfied if $m = 2\sin\vartheta\omega$. However, instead of capturing the resonance point directly, our approach is to take the average of the squared scattering amplitude over the possible incident angle uncertainty $\Delta\vartheta$ in QPS by the focused laser beams. Therefore, changing the focal length introduces different $\Delta\vartheta$, *i.e.*, the different range of the angular integral for the averaging. If a resonance peak is contained in that range, the resonance effect appears as the integrated result. The basic strategy of this proposal is thus to change the focal length, attempting to search for the appearance of four-wave mixing photons. The mass range is approximately given by the relation $m \le 2\omega\Delta\vartheta$. Therefore, this knob provides variations along the mass axis. On the other hand, the four-wave mixing yield is enhanced by the cubic product of the laser intensities. Therefore, changing laser intensities gives large variations along the coupling axis. If a significant signature is discovered, we are able to localize the domain in the mass-coupling relation by adjusting these two knobs.

Furthermore, there are knobs on the polarization states of laser fields. As we provided in Appendix of [8], depending on the types of exchanged fields, different polarization correlations are expected between the two photons in the initial and final states. In the case of the scalar field exchange which is the first of Eq. (1), the possible linear polarization states in the four-wave mixing process when all photons are on the same reaction plane, are expressed as follows:

$$\begin{aligned} \omega\{1\} + \omega\{1\} &\to (2-u)\omega\{1\} + u\omega\{1\} \\ \omega\{1\} + \omega\{1\} &\to (2-u)\omega\{2\} + u\omega\{2\}, \end{aligned} \tag{18}$$

where photon energies from the initial to the final state are denoted with the linear polarization states {1} and {2} which are orthogonal each other. On the other hand, in the case of the pseudoscalar field exchange with the second of Eq. (1), the possible linear polarization states are expressed as:

$$\begin{aligned} \omega\{1\} + \omega\{2\} &\to (2-u)\omega\{2\} + u\omega\{1\} \\ \omega\{1\} + \omega\{2\} &\to (2-u)\omega\{1\} + u\omega\{2\}. \end{aligned} \tag{19}$$

By choosing physically allowed combinations of the linear polarizations, we can distinguish the types of exchanged fields, while we can estimate the background processes by requiring the false combinations on purpose. This polarization dependence will be quite important to distinguish the atomic four-wave mixing process in the residual gas in the vacuum system in experiments. In order to understand backgrounds, we should gradually grade up the laser intensity.

As we discuss in the stimulation process, the key enhancement factor to improve the sensitivity is the number of photons or laser energy per pulse. On the other hand, if we consider many practical issues such as systematic understanding of background processes, we need a high repetition rate at the same time. We thus anticipate that laser facilities such as ICAN [15] can bring us the bright future for the dark field search.

In order to scan the broad mass range in the sub-eV domain, however, we further need to be able to introduce longer focal lengths, which can enhance the sensitivity to lower mass ranges via $\Delta\vartheta$, because a long focal length can introduce a small $\Delta\vartheta$. The focal length as a straight section would be limited up to order of 10 km in terrestrial experiments. On the other hand, for instance, if we could focus laser fields into the vacuum between two geostationary satellites, we expect to be able to use the distance of order of 10^4 km. We may transmit high-power laser beams into the one of the two satellites from the ground and focus them into the other one to perform the spectroscopy there. In such a case, a large telescope-type optics would be suitable to accurately handle the beam transmission.

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