Interpretation of excess in $H \to Z\gamma$ using a light axionlike particle

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We interpret the recent excess in a rare decay of the Higgs boson, $H \to Z\gamma$, using a light axionlike particle (ALP) in the mass range 0.05–0.1 GeV. The dominant decay of such a light ALP is into a pair of collimated photons, whose decay is required to happen before reaching the ECAL detector, such that it mimics a single photon in the detector. It can explain the excess with a coupling $C^{\rm eff}_{aZH}/\Lambda \sim 4\times 10^{-5}\ {\rm GeV}^{-1}$, while the decay of the ALP before reaching the ECAL requires the diphoton coupling $C^{\rm eff}_{\gamma\gamma}/\Lambda \geq 0.35\ {\rm TeV}^{-1}(0.1\ {\rm GeV}/m_a)^2$. A potential test would be the rare decay of the Z boson $Z \to aH^* \to a(b\bar{b})$ at the Tera-Z option of the future FCC and CEPC. However, it has a branching ratio of only $O(10^{-12})$, and thus barely testable. The production cross section for $pp \to Z^* \to aH$ via the same coupling $C^{\rm eff}_{aZH}/\Lambda$ at the LHC is too small for detection.

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I. INTRODUCTION

Since the discovery of the Higgs boson in 2012 [1,2], all the gauge couplings and the third-generation Yukawa couplings are shown to be consistent with the standard model (SM) Higgs boson (see the most recent fits [3]), including the loop-induced Hgg and $H\gamma\gamma$ couplings. The $H \to Z\gamma$ is one of the most anticipated measurements of the Higgs physics. Recently, an evidence of such a rare decay $H \to Z\gamma$ was jointly reported by ATLAS and CMS [4]. The search showed an observed significance of 3.4 standard deviations from the null hypothesis. The measured branching ratio of $H \to Z\gamma$,

$$B(H \to Z\gamma)_{\rm measured} = (3.4 \pm 1.1) \times 10^{-3}. \tag{1}$$

The SM prediction for the branching ratio of $H \to Z\gamma$ is [5]

$$B(H \to Z\gamma)_{\rm sm} = (1.5 \pm 0.1) \times 10^{-3}. \tag{2}$$

It is clear that the measurement showed an excess of 1.9σ [4].

In various well-founded extensions of the SM, there is a common occurrence of newly discovered pseudoscalar particles possessing masses lower than the electroweak scale. These particles serve various purposes, such as addressing the strong CP problem like axions [6–13] or acting as pseudoscalar mediators facilitating interaction between dark or hidden sectors and the SM [14]. Although it may be too early to conclude that new physics exists in the rare decay $H \to Z\gamma$, we speculate the interpretation of the excess using a very light axionlike particle (ALP) of mass $m_a = 0.05$–0.1 GeV. For such a light ALP the dominant decay of the ALP is $a \to \gamma\gamma$. Since the ALP is produced in the decay of the Higgs boson, we expect the transverse momentum p_{T_a} of the ALP is of order $(m_H/2)(1 - m_Z^2/m_H^2) \simeq m_H/4$, taking into account the massive Z boson. It is well known that the opening angle ΔR between the decay products of the ALP is then

$$\Delta R \sim \frac{2m_a}{p_{T_a}} \approx (3-7)\times 10^{-3},$$

for $m_a = 0.05$–0.1 GeV. We show the ΔR distributions for $m_a = 0.05$ GeV and 0.1 GeV in Fig. 1. Both the ATLAS and CMS detectors cannot resolve the two photons in such a small opening angle [15,16]. In this case, both photons deposit their energies in a single cell. In order that it happens, the axion has to decay before reaching or inside the ECAL detector. It is the coupling $C^{\rm eff}_{\gamma\gamma}/\Lambda$ in Eq. (5) that controls the decay length of the ALP. The ECAL detector the ATLAS detector extends from the radius of 1.5 to 2 m while that the CMS is slightly closer to the center. We therefore require the decay length of the axion to be less than 1.5 m. When these conditions are met, the diphoton

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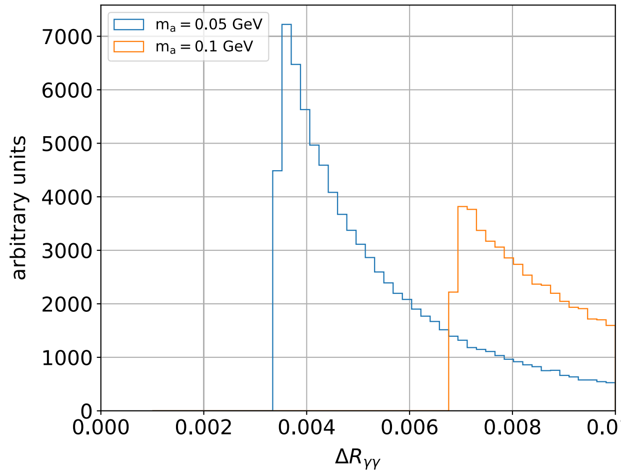

FIG. 1. Distributions of $\Delta R_{\gamma\gamma}$ between the photon pair produced for $m_a = 0.05$ GeV and 0.1 GeV in the decay $H \to Za \to (l^+l^-)(\gamma\gamma)$. It is clear the opening angle between the photon pair is very small.

decay of the axion would be mistaken as a single photon, and thus mimics the $H \to Z\gamma$ decay.

Details of the experimental event selections have been given in the CMS and ATLAS publications [17,18]. In both detectors, photons are identified as ECAL energy clusters not linked to the extrapolation of any charged particle trajectory to the ECAL. Typical angular resolution of the ECAL is of order $\Delta R \sim 10^{-2}$, which is given by the size of each cell, for photon energies of order 50–100 GeV [19,20]. It was demonstrated that taking advantage of a shower-shape analysis [21] (also emphasized in Ref. [22]) the ECAL detector is able to distinguish a single photon from a pair of collimated photons for $m_a \gtrsim 0.1$ GeV. That is the reason for the upper limit of $m_a = 0.1$ GeV that we propose while the lower limit 0.05 GeV is given by the existing limit on $C^{\rm eff}_{\gamma\gamma}/\Lambda$.

The final state of $Za \to (\ell^+\ell^-)(\gamma\gamma)$ mimics $(Z \to \ell^+\ell^-)\gamma$. Taking the difference between the measurement $B(H \to Z\gamma)_{\rm measured}$ and the SM prediction of $B(H \to Z\gamma)_{\rm sm}$ is entirely due to $H \to Za$, we obtain

$$B(H \to Za) = (1.9 \pm 1.1) \times 10^{-3}. \qquad (3)$$

In this work, we show that an effective coupling among a-Z-H, $C^{\rm eff}_{aZH}/\Lambda \approx 4.4 \times 10^{-5}\ {\rm GeV}^{-1}$, can explain the excess, without violating any other existing constraints. We also show that this interpretation may be tested at the Tera-Z option of the FCC [23] and CEPC [24]. On the other hand, the production cross section for $pp \to Z^* \to aH \to (\gamma\gamma)(b\bar{b})$ via the same coupling of aZH at the LHC is negligible for detection.

A few other interpretations were also proposed [25–27]. Barducci *et al.* [25] used extra chiral leptons with hypercharge Y and with scanning some choices of hypercharge Y the $H \to \gamma Z$ can be enhanced without increasing $H \to \gamma\gamma$. Boto *et al.* [26] used multiple charged scalar bosons S_i^{+Q}, which couple to H and Z, and with enhanced off-diagonal couplings, the authors can increase $H \to \gamma Z$ without increasing $H \to \gamma\gamma$. Das *et al.* [27] made use of the triplet scalar field in the context of Type II seesaw model and adjusted the couplings of the singly and doubly charged scalars to achieve the enhancement of $H \to Z\gamma$ without increasing $H \to \gamma\gamma$.

There are numerous studies of collimated photons from axion or ALP decays in literature; see for example [28–31].

II. MODEL

We follow the notation of Ref. [32]. The interactions of the ALP a with the SM particles start at dimension-5 [33],

$$\begin{aligned}\mathcal{L}^{D=5} &= \frac{1}{2}(\partial_\mu a)(\partial^\mu a) - \frac{1}{2}m_a^2 a^2 + \sum_f \frac{c_{ff}}{2\Lambda}\partial^\mu a \bar{f}\gamma_\mu\gamma_5 f \\ &\quad + g_S^2 \frac{C_{GG}}{\Lambda} a G^A_{\mu\nu}\tilde{G}^{\mu\nu,A} + g^2 \frac{C_{WW}}{\Lambda} a W^i_{\mu\nu}\tilde{W}^{\mu\nu,i} \\ &\quad + g'^2 \frac{C_{BB}}{\Lambda} a B_{\mu\nu}\tilde{B}^{\mu\nu},\end{aligned} \qquad (4)$$

where $A = 1, \ldots 8$ is the $SU(3)$ color index, $i = 1, 2, 3$ is the $SU(2)$ index, and g_S, g, and g' are the gauge couplings of $SU(3)$, $SU(2)$, and $U(1)_Y$, respectively. We set $C_{GG} = 0$ to avoid the mixing of the ALP with the QCD axion such that the strong CP problem would not come back. After the B_μ and W^3_μ rotate into the physical γ, Z, the ALP couples to γ and Z as

$$\begin{aligned}\mathcal{L} &= e^2 \frac{C_{\gamma\gamma}}{\Lambda} a F_{\mu\nu}\tilde{F}^{\mu\nu} + \frac{2e^2}{s_w c_w}\frac{C_{\gamma Z}}{\Lambda} a F_{\mu\nu}\tilde{Z}^{\mu\nu} \\ &\quad + \frac{e^2}{s_w^2 c_w^2}\frac{C_{ZZ}}{\Lambda} a Z_{\mu\nu}\tilde{Z}^{\mu\nu},\end{aligned} \qquad (5)$$

where

$$C_{\gamma\gamma} = C_{WW} + C_{BB}, \qquad C_{\gamma Z} = c_w^2 C_{WW} - s_w^2 C_{BB},$$
$$C_{ZZ} = c_w^4 C_{WW} + s_w^4 C_{BB},$$

and s_w and c_w are the sine and cosine of the weak mixing angle, respectively. In the considered mass range of the ALP $m_a \leq 0.1$ GeV, the only decay modes are e^+e^- and $\gamma\gamma$, for which the $\gamma\gamma$ can entirely dominate for $O(1)$ coefficients. However, we set $C_{ff} = 0$ for simplicity. Even in this case, the $a \to e^+e^-$ can be induced by a loop diagram, but it is largely suppressed. Therefore, the ALP so produced will decay entirely into a pair of photons, $\frac{C^{\rm eff}_{aZH}}{\Lambda}\frac{gv}{c_w}p^\mu$.

Interactions with the Higgs boson start at dimension-6,[1]

$$\mathcal{L}^{D\geq 6} = \frac{C_{ah}}{\Lambda^2}(\partial_\mu a)(\partial^\mu a)\phi^\dagger\phi + \frac{C_{aZH}}{\Lambda^3}(\partial^\mu a)\,(\phi^\dagger iD_\mu\phi + \text{H.c.})\phi^\dagger\phi, \quad (6)$$

where the covariant derivative is given by

$$D_\mu = \partial_\mu + i\frac{g}{\sqrt{2}}(W_\mu^+\tau^+ + W_\mu^-\tau^-) + ieQA_\mu + i\frac{g}{c_w}(T_3 - s_w^2 Q)Z_\mu,$$

and $\tau^\pm$ are the $SU(2)$ raising and lowering operators, T_3 is the third component of the isospin, and Q is the electric charge. It is easy to see that the first term in Eq. (6) induces the decay $H \to aa$ while the second term induces $H \to Za$. From dimensional analysis the amplitude for $H \to Za$ is suppressed by one more order of the cutoff scale Λ than $H \to aa$. However, as familiar to the Higgs low-energy theorems [35], in theories where a heavy new particle acquires most of its mass through electroweak symmetry breaking, the nonpolynomial dimension-5 operator can appear [22,32,36,37]

$$\frac{C_{aZH}^{(5)}}{\Lambda}(\partial^\mu a)(\phi^\dagger iD_\mu\phi + \text{H.c.})\ln(\phi^\dagger\phi/\mu^2), \quad (7)$$

which can be understood by thinking of ϕ as a background field and treating the heavy particle mass as a threshold for the running of gauge couplings.[2] Therefore, we can write an effective coupling for aZH as

$$C_{aZH}^{\text{eff}} = C_{aZH}^{(5)} + \frac{C_{aZH}v^2}{2\Lambda^2}. \quad (8)$$

We can now see that $H \to Za$ is only suppressed by one power of the cutoff scale Λ on amplitude level while $H \to aa$ by two powers of Λ. That is the reason why $H \to Za$ can be made sizable while keeping $H \to aa$ suppressed even in the case that both coefficients C_{ah} and C_{aZH}^{eff} are of order $O(1)$.

Note that the operator in Eq. (7) can induce a coupling among the H-a-$f\bar{f}$ after applying the equation of motion and integration by parts. Such a coupling can give rise to the rare decay $H \to b\bar{b}a$. Nevertheless, it is highly suppressed by Λ and the relatively small Yukawa m_b/v.

Before we end this section, we highlight existing constraints on other ALP-gauge couplings denoted by $g_{a\gamma\gamma}$, g_{aZZ}, $g_{aZ\gamma}$, and g_{aWW}. A dedicated study on the g_{aZZ}, $g_{aZ\gamma}$, and g_{aWW} was performed in Ref. [39] and references therein. These couplings can give rise to $pp \to Za \to (l^+l^-)(\gamma\gamma)$ and $pp \to Wa \to (l\nu)(\gamma\gamma)$, in which the photon pair can be resolved for larger m_a but unresolved for smaller m_a. Other existing collider constraints on g_{aZZ}, $g_{aZ\gamma}$, g_{aWW}, and $g_{a\gamma\gamma}$ have been discussed in Ref. [39]. Also, $Z \to a\gamma$, which looks like $Z \to \gamma\gamma$ when m_a is very small, was also searched for in $Z \to \gamma\gamma$ and summarized in [22]. On the other hand, comprehensive coverage of astrophysical constraints on $g_{a\gamma\gamma}$ can be found in [40].

III. RESULTS

For convenience of calculations we can write the effective vertex for aZH, after the electroweak symmetry breaking, as

$$\mathcal{L}_{aZH} = \frac{C_{aZH}^{\text{eff}}}{\Lambda}\frac{gv}{c_w}(\partial^\mu a)Z_\mu H, \quad (9)$$

which implies the Feynman rule in Fig. 2. Here $v \simeq 246$ GeV and c_w is the cosine of the Weinberg angle.

We can calculate the partial width of $H \to Za$ and $H \to aa$ [32]

$$\Gamma(H \to Za) = \frac{m_H^3}{16\pi}\left(\frac{C_{aZH}^{\text{eff}}}{\Lambda}\right)^2 \lambda^{3/2}(x_Z, x_a), \quad (10)$$

$$\Gamma(H \to aa) = \frac{m_H^3 v^2}{32\pi}\left(\frac{C_{aH}}{\Lambda^2}\right)^2 (1-2x_a)^2\sqrt{1-4x_a}, \quad (11)$$

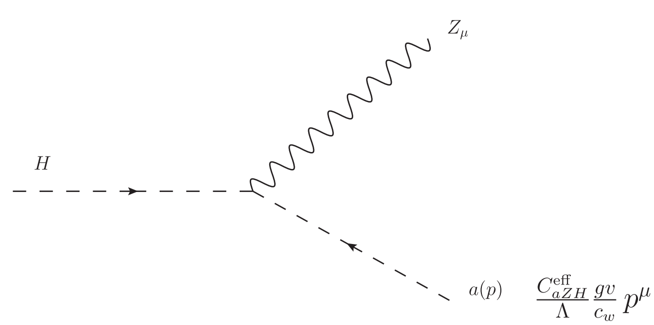

FIG. 2. Feynman rule for the vertex of aZH, which p_μ is the momentum of the incoming axion.

[1]The obvious dimension-5 operator $(\partial^\mu a)(\phi^\dagger iD_\mu\phi + \text{H.c.})$ is reduced to the fermionic operators using the equations of motion and integration by parts, such that this dimension-5 operator does not appear [34]. In other words, this dimension-5 operator cannot contribute to $H \to Za$ because there exists an equivalent basis in which this decay does not appear.

[2]A possibility of generating such an operator can be made by a triangular loop with a neutral heavy lepton N of mass TeV running in the loop, where N is the neutral component of an $SU(2)$ doublet and the charged component L is assumed to be much heavier. In such a setup, the $H\gamma\gamma$ and $Ha\gamma$ couplings are suppressed by the mass of L. Also, the HWW and HZZ couplings are unlikely to receive significant contributions from the triangular loops. The current mass limit on heavy neutral leptons is only about 600 GeV with $|V_{\mu N}|^2 \simeq 0.1$ [38].

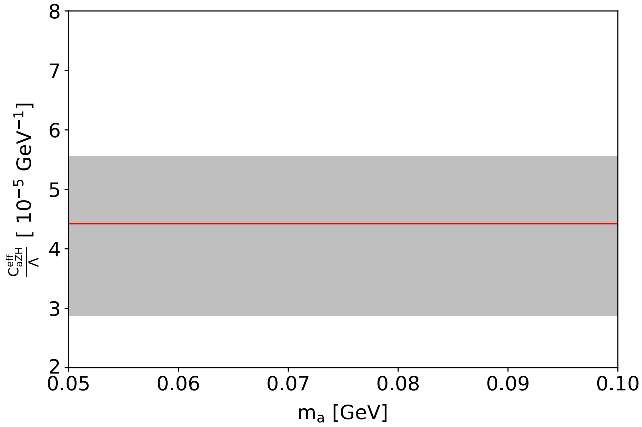

FIG. 3. The fitted values for $C^{\rm eff}_{aZH}/\Lambda$ versus m_a for $m_a = 0.05$–0.1 GeV. The red line and the band show the central value and 1σ uncertainty in $C^{\rm eff}_{aZH}/\Lambda = (4.4 \pm 1.1) \times 10^{-5}$ GeV^{-1} corresponding to $B(H \to Z\gamma) = (1.9 \times 1.1) \times 10^{-3}$.

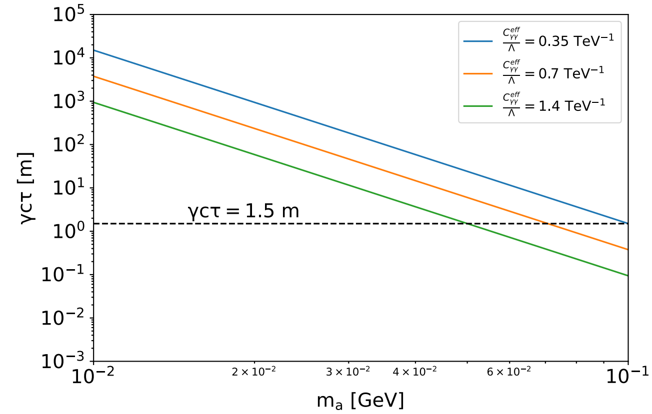

FIG. 4. Decay length $\gamma c\tau$ versus the mass m_a of the ALP. Values of $C^{\rm eff}_{\gamma\gamma}/\Lambda = 0.35, 0.7, 1.4$ TeV^{-1} are used. A dashed horizontal line of 1.5 m is also shown.

where $x_i = m_i^2/m_H^2 (i = a, Z)$ and $\lambda(x, y) = (1 - x - y)^2 - 4xy$. Including the new contribution of $\Gamma(H \to Za)$ the branching ratio of $H \to Za$ is given by

$$B(H \to Za) = \frac{\Gamma(H \to Za)}{\Gamma(H \to Za) + \Gamma_{\rm sm}(m_H = 125\,{\rm GeV})}, \tag{12}$$

where $\Gamma_{\rm sm}(m_H = 125\ {\rm GeV})$ is taken to be 4.088×10^{-3} GeV for $m_H = 125$ GeV.[3]

Requiring the branching ratio to be $(1.9 \pm 1.1) \times 10^{-3}$ as in Eq. (3), we obtain the results as shown in Fig. 3. We find that

$$\frac{C^{\rm eff}_{aZH}}{\Lambda} = (4.4^{+1.1}_{-1.6}) \times 10^{-5}\ {\rm GeV}^{-1}, \tag{13}$$

where the upper and lower limits correspond to the 1σ of $B(H \to Za) = (1.9 \pm 1.1) \times 10^{-3}$. If the coefficient $C^{\rm eff}_{aZH} \sim O(1)$ the corresponding cutoff scale is $\Lambda = 22.6$ TeV.

It is true that the result of $C^{\rm eff}_{aZH}/\Lambda$ corresponds to the mass scale of $\Lambda = 22.6$ TeV with $O(1)$ coefficient. If it is the case, these heavy particles would certainly be out of reach at the LHC. On the other hand, if we take the coefficient to be $O(0.1) \approx e^2$ [as in the definition of $e^2 (C_{\gamma\gamma}/\Lambda) a F_{\mu\nu} \tilde{F}^{\mu\nu}$], the scale Λ would then be around 2 TeV, which may be readily available at the LHC. Indeed, the current mass limit on heavy vectorlike quarks is about $O(1) - 1.6$ TeV depending on the search channels (for a recent review see [41]), and the mass limit on heavy neutral leptons is about 600 GeV with $|V_{\mu N}|^2 \simeq 0.1$ [38].

Let us turn to the requirement on the coupling $C^{\rm eff}_{\gamma\gamma}/\Lambda$, which controls the decay length $\gamma c\tau$ of the ALP, where $\gamma = E_a/m_a (E_a \approx m_H/2)$ is the Lorentz boost factor of the ALP and the decay time τ in the rest frame is given by

$$\tau = \frac{1}{\Gamma_a}, \qquad \Gamma_a = 4\pi\alpha^2 m_a^3 \left(\frac{C^{\rm eff}_{\gamma\gamma}}{\Lambda}\right)^2, \tag{14}$$

where Γ_a is the total decay width of the ALP assuming it only decays into diphoton. We show the decay length of the ALP versus m_a for a few values of $C^{\rm eff}_{\gamma\gamma}/\Lambda$ in Fig. 4. Taking the input values of $E_a = m_H/2 = 62.5$ GeV, the requirement of the $\gamma c\tau \leq 1.5$ m gives

$$\frac{C^{\rm eff}_{\gamma\gamma}}{\Lambda} \geq 0.35\ {\rm TeV}^{-1} \left(\frac{0.1\ {\rm GeV}}{m_a}\right)^2. \tag{15}$$

Therefore, at $m_a = 0.1 (0.05)$ GeV the coupling $C^{\rm eff}_{\gamma\gamma}/\Lambda > 0.35(1.4)$ TeV^{-1}. We show in Fig. 5 the region of parameter space in $(m_a, C^{\rm eff}_{\gamma\gamma}/\Lambda)$ that can allow the ALP to decay before reaching the ECAL and consistent with all existing constraints. Note that the lower mass limit $m_a = 0.05$ GeV is due to the existing constraints (see Fig. 5), while the upper limit $m_a = 0.1$ GeV came from the limitation of the shower-shape analysis [21].

Such a scenario using a light axion with the diphoton decay, which mimics a single photon, to explain the excess in $H \to Z\gamma$ can be tested at the Z resonance (Tera-Z – $10^{12} Z$ bosons) of the Future Circular Colliders [23] and CEPC [24]. Via the same coupling $C^{\rm eff}_{aZH}/\Lambda$ the Z boson can decay via an off shell Higgs boson,

$$Z \to aH^* \to a(b\bar{b}),$$

in which the most dominant mode of the virtual Higgs boson is considered. The final state consists of a $b\bar{b}$ pair plus a diphoton, which appears as a single photon.

[3] It is available at https://twiki.cern.ch/twiki/bin/view/LHCPhysics/CERNYellowReportPageBR.

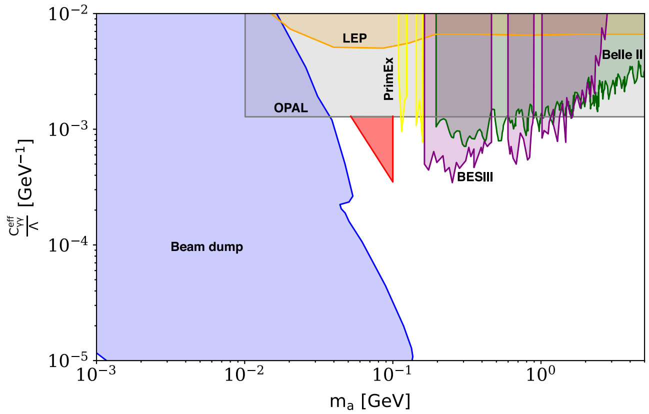

FIG. 5. Parameter space (shaded in red) in $C^{\rm eff}_{\gamma\gamma}/\Lambda$ versus m_a that can allow the ALP to decay before reaching the ECAL (i.e., $\gamma c\tau \le 1.5$ m) and consistent with all existing constraints in the mass range of $10^{-3} - 5$ GeV, including beam dump [42–46], OPAL [47], LEP [48], Belle II [49], BES III [50], and PrimEx [51] (data extracted from the GitHub page [52]). Note that the mass range of the fitted parameter space is 0.05 GeV $\le m_a \le 0.1$ GeV.

Nevertheless, the branching ratio is only 10^{-12}, which barely affords a few events at the Tera-Z option.

Another possible test of the scenario is the production process $pp \to Z^* \to aH$ at the LHC via the same coupling $C^{\rm eff}_{aZH}/\Lambda$. However, the cross section turns out to be negligible, of order 10^{-6} fb only, which corresponds to far less than one event for 3000 fb^{-1} luminosity of the entire running of high-luminosity LHC (HL-LHC).

IV. CONCLUSIONS

The excess observed in the rare decay of the Higgs boson into a Z boson and a photon can be interpreted as the Higgs decay into a Z boson and a light axion. The light axion then decays into a pair of collimated photons such that the ECAL cannot resolve. Such a scenario requires a coupling between aZH with strength $C^{\rm eff}_{aZH}/\Lambda \sim 4\times 10^{-5}$ GeV^{-1}. Furthermore, the axion is required to decay before reaching the ECAL, which implies the effective $C^{\rm eff}_{\gamma\gamma}/\Lambda \ge 0.35$ TeV$^{-1}(0.1\ \text{GeV}/m_a)^2$. Such a aZH coupling may be tested via $Z \to aH^* \to a(b\bar{b})$ at the Tera-Z option of the FCC and CEPC, but, however, it has a branching ratio of only 10^{-12}.

Note that the cutoff scale defined in $C^{\rm eff}_{aZH}/\Lambda$ is about 22.6 TeV with $O(1)$ coefficient. On the hand, the requirement of the decay length of $\gamma c\tau \le 1.5$ m needs $C^{\rm eff}_{\gamma\gamma}/\Lambda = 0.35$ TeV$^{-1}(0.1\ \text{GeV}/m_a)^2$. Note that there is a factor $e^2 \sim 0.1$ in front of $C_{\gamma\gamma}/\Lambda$ in Eq. (5). Therefore, if we take out this factor e^2, the corresponding Λ with $O(1)$ coefficient would become 28 TeV for $m_a = 0.1$ GeV and 7 TeV for $m_a = 0.05$ GeV, such that these two sets of Λ's are of similar order. Nevertheless, these two values are purely phenomenological.

A comment on the constraints from flavor-changing processes is in order here. The ALP with mass $m_a =$ 0.05–0.1 GeV may subject to constraints from flavor-changing processes such as $K \to \pi\nu\bar{\nu}$ and $K \to \pi\mu^+\mu^-$ and the corresponding ones for $B \to K$. It was shown in Ref. [53] that C_{WW} coupling can induce ALP flavor-changing coupling at one-loop order such that the constraints on $C^{\rm eff}_{\gamma\gamma}/\Lambda$ strongly restrict our fitted region of Fig. 4. On the other hand, if only C_{BB} coupling exists in the UV scale, the flavor-changing couplings involving the ALP can only be generated at two loops [53], and are therefore highly suppressed. As shown in the right panel of Fig. 22 of Ref. [53], the flavor constraints are rather weak in this case and our fitted parameter space is valid.

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