Diphoton signatures from heavy axion decays at the CERN Large Hadron Collider

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Recently, the LHC collaborations, ATLAS and CMS, have announced an excess in the diphoton channel with local significance of about 3σ around an invariant mass distribution of ∼750 GeV, after analyzing new data collected at center-of-mass energies of $\sqrt{s} = 13$ TeV. We present a possible physical interpretation of such a signature, within the framework of a minimal UV-complete model with a massive singlet pseudoscalar state a that couples to a new TeV-scale colored vectorlike fermion F, whose hypercharge quantum number is a non-zero integer. The pseudo-scalar state a might be a heavy pseudo-Goldstone boson, such as a heavy axion, which decays into two photons and whose mass lies around the excess region. The mass of the CP-odd state a and its coupling to F may be due to nonperturbative effects, which can break the original Goldstone shift symmetry dynamically. The possible role that the heavy axion a can play in the radiative generation of a seesaw Majorana scale and in the solution to the so-called strong CP problem is briefly discussed.

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Recently, the ATLAS and CMS collaborations have analyzed Run 2 LHC data gathered at center-of-mass energies of $\sqrt{s} = 13$ TeV. They reported an excess in the diphoton channel around an invariant mass distribution of ∼750 GeV, with local significance of 3.6σ and 2.6σ confidence level (CL), respectively [1]. Interestingly enough, Run 2 data do not show up any significant excess in other diboson channels, such as ZZ, W^+W^- and $Z\gamma$, whilst the Run–1 bumps seen around the 2 TeV region have now become almost statistically insignificant.

In this short paper, we offer a possible interpretation of such an excess in the diphoton channel, within the framework of a minimal UV-complete model with a massive singlet pseudoscalar state a that couples to new colored vectorlike fermions $F_{L,R}$. These vectorlike fermions are very heavy with masses $m_F \gtrsim 1.5$ TeV, so as to have escaped detection so far at the LHC. They are charged under the $\mathrm{SU}(3)_C$ group of quantum chromodynamics (QCD), but they are singlets under the weak $\mathrm{SU}(2)_L$ group of the standard model (SM). They must also have nonzero integer hypercharges, e.g., $Y_{F_L} = Y_{F_R} = 1, 2, \ldots,$ which forbid them to have Yukawa interactions with the SM quarks. On the other hand, the pseudoscalar state a may well be a heavy pseudo-Goldstone boson, such as a heavy axion, which decays into two photons with a mass that lies around the excess region of ∼750 GeV. Both the mass M_a of the state a and its CP-odd coupling to the Dirac vectorlike fermion F, $h_F \bar{F} i\gamma_5 F$, could originate from nonperturbative effects that break the original axion shift symmetry.

The minimal UV-complete model that we will be considering here is related to the one that discussed earlier in Ref. [2]. The relevant non-SM part of the Lagrangian of interest to us is given by

$$\mathcal{L} = \bar{F}(i\not{D} - m_F)F + \frac{1}{2}(\partial_\mu a)(\partial^\mu a) - \frac{1}{2}M_a^2 a^2 - h_F a \bar{F} i\gamma_5 F. \tag{1}$$

In the above, $D_\mu = \partial_\mu + i g_s T^a G_\mu^a + i g'(Y_F/2) B_\mu$ is the covariant derivative acting on the exotic colored Dirac fermion F, where G_μ^a and B_μ are the $\mathrm{SU}(3)_C$ and $\mathrm{U}(1)_Y$ gauge bosons, respectively, and T^a (with $a = 1, 2, \ldots, 8$) are the generators of the $\mathrm{SU}(3)_C$ gauge group. Notice that Lagrangian (1) is invariant under the CP transformations: $a(t, \mathbf{x}) \to -a(t, -\mathbf{x})$ and $\bar{F}(t, \mathbf{x}) i\gamma_5 F(t, \mathbf{x}) \to -\bar{F}(t, -\mathbf{x}) i\gamma_5 F(t, -\mathbf{x})$. In the absence of the fermion mass term $m_F \bar{F} F$, Lagrangian (1) is also invariant under the chirality discrete transformations: $a \to -a$ and $F_{R(L)} \to +(-)F_{R(L)}$. Given that $m_F \neq 0$, this latter symmetry is broken softly by the dimension-3 mass operator $m_F \bar{F} F$. Finally, it is important to remark that the squared mass M_a^2 and the Yukawa couplings h_F in Lagrangian (1) break explicitly the Goldstone shift symmetry: $a \to a + c$, where c is an arbitrary constant. The possible origin of such a breaking could be due to nonperturbative effects related to some unspecified strong dynamics.

In the above minimal extension of the SM, the pseudoscalar field a couples to the electromagnetic (em) field A_μ and the gluon fields G_μ^a, via the five-dimensional operators: $aF^{\mu\nu}\tilde{F}_{\mu\nu}$ and $aG^{a\mu\nu}\tilde{G}_{\mu\nu}^a$, where $F^{\mu\nu}$ and $G^{a,\mu\nu}$ are the $\mathrm{U}(1)_{\mathrm{em}}$ and $\mathrm{SU}(3)_C$ field strength tensors, respectively, and $\tilde{F}_{\mu\nu} \equiv \frac{1}{2}\varepsilon_{\mu\nu\lambda\rho}F^{\lambda\rho}$ and $\tilde{G}_{\mu\nu}^a \equiv \frac{1}{2}\varepsilon_{\mu\nu\lambda\rho}G^{a,\lambda\rho}$ are their corresponding dual tensors. These operators are induced by the chiral global anomalies of the heavy fermion F, through the triangle graphs shown in Fig. 1. With the convention that all

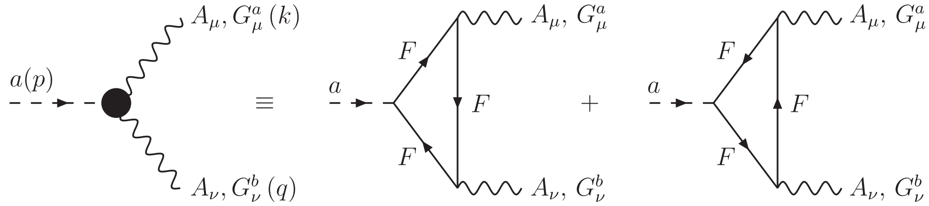

FIG. 1. The operators $aF_{\mu\nu}\tilde{F}^{\mu\nu}$ and $aG^a_{\mu\nu}\tilde{G}^a_{\mu\nu}$, as induced by the chiral global anomaly of the heavy fermion F, with the convention $p+k+q=0$.

momenta are incoming, i.e., $p+k+q=0$, the one-loop $a(p)-A_\mu(k)-A_\nu(q)$ coupling reads [3–5]:

$$i\Gamma^{aAA}_{\mu\nu}(p,k,q) = iQ_F^2\frac{N_C\alpha_{\rm em}}{\pi}\frac{h_F}{m_F}F_P\left(\frac{p^2}{4m_F^2}\right)\varepsilon_{\mu\nu\lambda\rho}k^\lambda q^\rho, \qquad (2)$$

where $Q_F = Y_F/2$ is the electric charge of the heavy fermion F, $N_C = 3$ is its color degrees of freedom, $\alpha_{\rm em} = e^2/(4\pi)$ is the electromagnetic fine structure constant, and $\varepsilon_{\mu\nu\lambda\rho}$ is the usual antisymmetric Levi–Civita tensor (with the convention: $\varepsilon^{0123} = +1$). Moreover, the loop function $F_P(\tau)$ was calculated a long time ago [3] and found to be

$$F_P(\tau) = \begin{cases} \frac{1}{\tau}\arcsin^2\sqrt{\tau}; & |\tau|\le 1, \\ -\frac{1}{4\tau}\left[\ln\left(\frac{\sqrt{\tau}+\sqrt{\tau-1}}{\sqrt{\tau}-\sqrt{\tau-1}}\right) - i\pi\right]^2; & |\tau|\ge 1. \end{cases} \qquad (3)$$

Note that for $|\tau| \ll 1$, we have $F_P(\tau) = 1 + \tau/3 + \mathcal{O}(\tau^2)$, whereas for $|\tau| \gg 1$, $F_P(\tau) \to -\ln^2|\tau|/(4\tau)$ which goes to zero asymptotically as $\tau\to\infty$.

By analogy, the $\mathrm{SU}(3)_C$ global anomaly generates an effective interaction of the heavy axion a to gluons G^a_μ, as shown in Fig. 1. The effective $a(p)-G^a_\mu(k)-G^b_\nu(q)$ coupling is given by

$$i\Gamma^{aG^aG^b}_{\mu\nu}(p,k,q) = i\delta^{ab}\frac{\alpha_s}{2\pi}\frac{h_F}{m_F}F_P\left(\frac{p^2}{4m_F^2}\right)\varepsilon_{\mu\nu\lambda\rho}k^\lambda q^\rho, \qquad (4)$$

where $\alpha_s = g_s^2/(4\pi)$ is the strong fine structure constant.

With the aid of the effective couplings given in (2) and (4), it is straightforward to calculate the decay widths of the heavy axion a into photons (γ) and gluons (g):

$$\Gamma(a\to\gamma\gamma) = \frac{N_C^2\alpha_{\rm em}^2}{64\pi^3}Q_F^4 h_F^2\frac{M_a^3}{m_F^2}|F_P(\tau_a)|^2, \qquad (5)$$

$$\Gamma(a\to gg) = \frac{\alpha_s^2}{32\pi^3}h_F^2\frac{M_a^3}{m_F^2}|F_P(\tau_a)|^2 K_a^g, \qquad (6)$$

where $\tau_a \equiv M_a^2/(4m_F^2)$ and $K_a^g \approx 1.6$ is a QCD loop enhancement factor which includes the leading order QCD corrections [6]. In addition, the other diboson decay channels, such as $a\to ZZ$, $Z\gamma$ and W^+W^-, may be reliably estimated to leading order in M_Z^2/M_a^2 [7] as follows:

$$\frac{\Gamma(a\to ZZ)}{\Gamma(a\to\gamma\gamma)} \approx \frac{\sin^4\theta_w}{\cos^4\theta_w} \approx 0.082, \qquad \frac{\Gamma(a\to Z\gamma)}{\Gamma(a\to\gamma\gamma)} \approx \frac{2\sin^2\theta_w}{\cos^2\theta_w} \approx 0.57, \qquad (7)$$

whilst the decay width $a\to W^+W^-$ is negligible, since the corresponding aW^+W^- effective coupling is generated at the two-loop level, e.g., from the one-loop induced $a\gamma\gamma$ coupling. To satisfy the LHC constraints on the masses of exotic colored fermions, we may assume that the vectorlike fermion F is heavier than a, e.g., $m_F \gtrsim 1.5$ TeV, in which case $\tau_a \ll 1$. Hence, the loop function $F_P(\tau_a)$ may well be approximated as $F_P(\tau_a)\approx 1$.

If we now take the ratio R of the photonic versus the gluonic decay width given in (5) and (6), we readily find that

$$R \equiv \frac{\Gamma(a\to\gamma\gamma)}{\Gamma(a\to gg)} = \frac{N_C^2\alpha_{\rm em}^2 Q_F^4}{2\alpha_s^2 K_a^g}. \qquad (8)$$

Observe that the ratio R is independent of the Yukawa coupling h_F and, for $Q_F\ge 2$, we obtain $R>1$ and the decay $a\to\gamma\gamma$ can easily become the dominant mode.

The production cross section of heavy axions via gluon-gluon fusion [8] may be calculated as follows:

$$\sigma(pp\to a\to\gamma\gamma) \approx \sigma(pp\to a)B(a\to\gamma\gamma), \qquad (9)$$

where $B(a\to\gamma\gamma)\approx R/(1+1.57R)$ is the branching fraction for the decay channel $a\to\gamma\gamma$, with R given in (8). For center-of-mass energies of $\sqrt{s}=13$ TeV, we may naively estimate the cross section $\sigma(pp\to a)$ as

$$\sigma(pp\to a) \sim \sigma_{\rm SM}(pp\to H)\times h_F^2\frac{m_t^2}{m_F^2}\frac{M_a^2}{M_H^2}, \qquad (10)$$

where $\sigma_{\rm SM}(pp\to H)\approx 40$ pb is a reference production cross section of the SM Higgs boson H via gluon-gluon fusion, with $M_H\approx 125$ GeV [9]. Hence, for $M_a = 750$ GeV (or $M_a/M_H = 6$), $m_F/m_t = 10$ and $h_F = 0.1$, we find that

$$\sigma(pp\to a\to\gamma\gamma) \sim 15\ \text{fb}\times B(a\to\gamma\gamma). \qquad (11)$$

For $B(a \to \gamma\gamma) \sim 1$ and an integrated luminosity $\mathcal{L} = 3\ \mathrm{fb}^{-1}$ at $\sqrt{s} = 13$ TeV, we obtain about 45 signal events, which is compatible with the diphoton-excess events reported in [1].

As discussed in detail in [2], axionlike fields could act as mediators to generate TeV-scale gauge-invariant masses, such as $m_F \bar{F}F$, for vectorlike fermions through global anomalies at the three-loop level. In particular, a gauge-invariant Majorana mass term $m_M (\bar{\nu}_R)^C \nu_R$ can be generated [10], if heavy axion fields couple to Kalb–Ramond axions [11,12] that occur in torsionful theories of quantum gravity. Light axions play an important role in solving the strong CP problem via the so-called Peccei–Quinn mechanism [13–16]. Thus, the possible presence of a heavy axion, or a multitude of axions [17,18], may give rise to interesting mixing phenomena and possibly to new effects in astrophysical considerations [19].

In conclusion, we have presented a minimal UV-complete model, based on the possible existence of a heavy axion with mass $M_a \approx 750$ GeV, which could offer a possible physical interpretation of the diphoton excess observed in the Run 2 data. The model requires the presence of a new TeV-scale colored vectorlike fermion F, which has a nonzero integer hypercharge. For large electric charges $Q_F \geq 2$, the photonic decay mode $a \to \gamma\gamma$ becomes naturally the dominant channel. The latter, along with the branching fractions given in (7), provide a unique prediction of our minimal model that can be tested with future Run 2 data. Nevertheless, our model may require an extension to its field content, as it exhibits a Landau pole at energy scales $Q \lesssim 10^{14}$ GeV, for $Q_F \geq 2$. Further studies are therefore needed, so as to be able to fully assess the physical significance of the observed diphoton excess, as a firm signature of new physics at the LHC.

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