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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 $SU(3)_C$ group of quantum chromodynamics (QCD), but they are singlets under the weak $SU(2)_L$ group of the standard model (SM). They must also have nonzero integer hypercharges, e.g., $Y_{F_L} = Y_{F_R} = 1, 2, \dots$, 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

$$\mathcal{L} = \bar{F}(iD - m_F)F + \frac{1}{2}(\partial_\mu a)(\partial^\mu a) - \frac{1}{2}M_a^2 a^2 - h_F a \bar{F} i \gamma_5 F. \quad (1)$$

In the above, $D_\mu = \partial_\mu + ig_s T^a G_\mu^a + ig'(Y_F/2)B_\mu$ is the covariant derivative acting on the exotic colored Dirac fermion F , where G_μ^a and B_μ are the $SU(3)_C$ and $U(1)_Y$ gauge bosons, respectively, and T^a (with $a = 1, 2, \dots, 8$) are the generators of the $SU(3)_C$ gauge group. Notice that Lagrangian (1) is invariant under the CP transformations: $a(t, \mathbf{x}) \rightarrow -a(t, -\mathbf{x})$ and $\bar{F}(t, \mathbf{x}) i \gamma_5 F(t, \mathbf{x}) \rightarrow -\bar{F}(t, -\mathbf{x}) i \gamma_5 F(t, -\mathbf{x})$. In the absence of the fermion mass term $m_F \bar{F} F$, Lagrangian (1) is also invariant under the chirality discrete transformations: $a \rightarrow -a$ and $F_{R(L)} \rightarrow +(-)F_{R(L)}$. 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: $a \rightarrow a + c$, 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 pseudo-scalar field a couples to the electromagnetic (em) field A_μ and the gluon fields G_μ^a , via the five-dimensional operators: $a F^{\mu\nu} \tilde{F}_{\mu\nu}$ and $a G^{a\mu\nu} \tilde{G}_{\mu\nu}^a$, where $F^{\mu\nu}$ and $G^{a,\mu\nu}$ are the $U(1)_{\text{em}}$ and $SU(3)_C$ field strength tensors, respectively, and $\tilde{F}_{\mu\nu} \equiv \frac{1}{2} \epsilon_{\mu\nu\lambda\rho} F^{\lambda\rho}$ and $\tilde{G}_{\mu\nu}^a \equiv \frac{1}{2} \epsilon_{\mu\nu\lambda\rho} G^{a,\lambda\rho}$ 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 $aG_{\mu\nu}^a\tilde{G}_{\mu\nu}^a$, 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]:

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

where $Q_F = Y_F/2$ is the electric charge of the heavy fermion F , $N_C = 3$ is its color degrees of freedom, $\alpha_{\text{em}} = e^2/(4\pi)$ 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

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

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

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

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

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):

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

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

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 $a \rightarrow ZZ$, $Z\gamma$ and W^+W^- , may be reliably estimated to leading order in M_Z^2/M_a^2 [7] as follows:

$$\begin{aligned} \frac{\Gamma(a \rightarrow ZZ)}{\Gamma(a \rightarrow \gamma\gamma)} &\approx \frac{\sin^4 \theta_w}{\cos^4 \theta_w} \approx 0.082, \\ \frac{\Gamma(a \rightarrow Z\gamma)}{\Gamma(a \rightarrow \gamma\gamma)} &\approx \frac{2\sin^2 \theta_w}{\cos^2 \theta_w} \approx 0.57, \end{aligned} \quad (7)$$

whilst the decay width $a \rightarrow W^+W^-$ 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

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

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

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

$$\sigma(pp \rightarrow a \rightarrow \gamma\gamma) \approx \sigma(pp \rightarrow a) B(a \rightarrow \gamma\gamma), \quad (9)$$

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

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

where $\sigma_{\text{SM}}(pp \rightarrow H) \approx 40$ 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

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

For $B(a \rightarrow \gamma\gamma) \sim 1$ and an integrated luminosity $\mathcal{L} = 3 \text{ fb}^{-1}$ at $\sqrt{s} = 13 \text{ TeV}$, 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 $M_a \approx 750 \text{ GeV}$, 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 $a \rightarrow \gamma\gamma$ 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 $Q \lesssim 10^{14} \text{ GeV}$, 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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