



Diphoton excess and running couplings



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ARTICLE INFO

Article history:

Received 22 February 2016

Accepted 14 April 2016

Available online 18 April 2016

Editor: J. Hisano

ABSTRACT

The recently observed diphoton excess at the LHC may suggest the existence of a singlet (pseudo-)scalar particle with a mass of 750 GeV which couples to gluons and photons. Assuming that the couplings to gluons and photons originate from loops of fermions and/or scalars charged under the Standard Model gauge groups, we show that there is a model-independent upper bound on the cross section $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ as a function of the cutoff scale Λ and masses of the fermions and scalars in the loop. Such a bound comes from the fact that the contribution of each particle to the diphoton event amplitude is proportional to its contribution to the one-loop β functions of the gauge couplings. We also investigate the perturbativity of running Yukawa couplings in models with fermion loops, and show the upper bounds on $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ for explicit models.

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1. Introduction

Recently, the ATLAS and CMS Collaborations reported an excess of diphoton events implying a resonance with a mass of around 750 GeV [1,2]. The ATLAS Collaboration has 3.2 fb^{-1} of data, and the largest excess is found at around the diphoton invariant mass of $m_{\gamma\gamma} \simeq 750 \text{ GeV}$ with the local (global) significance of 3.6σ (2.0σ) for a narrow width case. When a large width for the signal component is assumed, the local (global) significance increases to 3.9σ (2.3σ) at the width of about 45 GeV. The CMS Collaboration, with 2.6 fb^{-1} of data, also reported an excess at around $m_{\gamma\gamma} \simeq 750 \text{ GeV}$ with the local (global) significance of 2.6σ (1.2σ) for a narrow width case, while the significance does not increase with a larger width. Possible explanations and implications of this excess have been extensively discussed [3–7].

One of the plausible explanations of the excess is that a scalar or pseudoscalar particle S with a mass of 750 GeV is produced through gluon fusion and decays into a pair of photons, $gg \rightarrow S \rightarrow \gamma\gamma$, via diagrams with new fermions and/or bosons charged under the Standard Model (SM) gauge groups running in the loops [3–6]. In order to explain the excess with perturbative couplings, however, the new particles in the loop should have large quantum numbers and/or large multiplicity, which implies that the perturbativity of the SM gauge groups may break down at some high

scale below the Planck scale. In this letter, we address this issue and investigate the perturbativity of such models.

Our main conclusions are as follows:

1. We point out that the contribution of each particle in the loop to the diphoton event amplitude is proportional to its contribution to the one-loop β functions of the gauge couplings at the leading order, independently of the representations of the particles in the loop. Consequently, there is a generic upper bound on the cross section $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ as a function of the cutoff scale Λ and masses of the fermions and scalars in the loop. We also numerically evaluate such a bound, taking into account the following constraints:
 - (i) the constraints from Landau pole, requiring that the gauge couplings remain perturbative up to the scale Λ , and
 - (ii) the constraint from the scale dependence of the strong coupling constant based on the LHC [8].
2. We also investigate the running of the Yukawa coupling in models with fermion loops. The upper bound on $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ is presented as a function of the fermion mass and the cutoff scale Λ for some explicit models with vector-like quarks.

The generic analysis in the first part, which can be applied to models with fermions and scalars in the loop in arbitrary representations, was not considered in the previous works. The analysis of the second part is close to those of Ref. [3], where the authors

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investigated the running of the gauge, Yukawa, and scalar quartic couplings in models with multiple generations of fermions in the loop. (See also Refs. [4] for related works.) They considered several model points with fixed fermion masses and the number of generations. Our analysis is complementary in the sense that the fermion mass, the Yukawa coupling, the number of generations, as well as the cutoff scale are taken as free parameters.

In the next section, we investigate the running of gauge couplings in generic setup with fermions and scalars in arbitrary representations, and show that there is a model-independent upper bound on the cross section $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$. In Sec. 3, we investigate the running of the Yukawa coupling (as well as those of gauge couplings) in explicit models and present the upper bound on $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ as a function of the fermion mass and the cutoff scale Λ . We also briefly discuss the LHC constraints on vector-like quarks, and comment on the running of the scalar quartic coupling of S . We conclude in Sec. 4.

2. Running gauge couplings and generic upper bound on the diphoton event rate

The reported diphoton excess can be explained by a new scalar particle S , with a mass of $m_S \simeq 750$ GeV, which is produced by a gluon fusion and decays into two photons. The cross section is given by

$$\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = \frac{C_{gg}}{s \cdot m_S} \frac{\Gamma(S \rightarrow gg)\Gamma(S \rightarrow \gamma\gamma)}{\Gamma_{S,\text{total}}}, \quad (1)$$

where $\sqrt{s} = 13$ TeV is the center-of-mass energy of the LHC, and $C_{gg} = (\pi^2/8) \int_0^1 dx_1 \int_0^1 dx_2 \delta(x_1 x_2 - m_S^2/s) g(x_1) g(x_2)$, with $g(x)$ being the gluon parton distribution function. In our numerical calculation, we use the MSTW2008 NLO set [9] evaluated at the scale $\mu = m_S$, which gives $C_{gg} \simeq 2.1 \times 10^3$. The reported excess [1,2] suggests $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) \sim \mathcal{O}(1)\text{--}10$ fb.

We assume that the production and the decay of the singlet scalar S is induced through loops of new fermions ψ_i and/or scalars ϕ_i . In order to make the analysis model-independent, we consider that they have generic quantum numbers $(R_i^{(3)}, R_i^{(2)}, Y_i)$ under the SM gauge groups $SU(3) \times SU(2) \times U(1)_Y$. The relevant part of the Lagrangian is given by¹

$$\begin{aligned} \mathcal{L} = & \mathcal{L}_{\text{SM}} + \frac{1}{2}(\partial_\mu S)^2 - \frac{1}{2}m_S^2 S^2 \\ & + \sum_{i=\text{fermions}} \eta_i (\bar{\psi}_i(i\not{\partial} - m_i)\psi_i - y_i S \bar{\psi}_i \psi_i - iy_{5i} S \bar{\psi}_i \gamma_5 \psi_i) \\ & + \sum_{i=\text{scalars}} \eta_i (|D_\mu \phi_i|^2 - m_i^2 |\phi_i|^2 - A_i S |\phi_i|^2) \\ & + (\text{scalar quartic couplings}), \end{aligned} \quad (2)$$

where $\eta_i = 1/2$ for Majorana fermions and real scalars, and $\eta_i = 1$ otherwise. (Notice that Majorana fermions and real scalars are possible only for the case of real representation of the SM gauge group, such as $(\mathbf{8}, \mathbf{1}, 0)$ and $(\mathbf{1}, \mathbf{3}, 0)$.)

In the following, we assume CP-conservation, and consider the two cases of scalar S ($y_{5i} = 0$) and pseudoscalar S ($y_i = A_i = 0$) separately. The partial decay rates of S into gg and $\gamma\gamma$ are given by

$$\Gamma(S \rightarrow gg) = \frac{2}{\pi} \kappa_{gg}^2 m_S^3, \quad \Gamma(S \rightarrow \gamma\gamma) = \frac{1}{4\pi} \kappa_{\gamma\gamma}^2 m_S^3, \quad (3)$$

¹ In general, off-diagonal couplings such as $y_{ij} S \bar{\psi}_i \psi_j$ and/or $A_{ij} S \phi_i^* \phi_j$ ($i \neq j$) are allowed when ψ_i and ψ_j (ϕ_i and ϕ_j) have the same quantum numbers, but they do not contribute to the process $gg \rightarrow S \rightarrow \gamma\gamma$ at the one loop.

where

$$\begin{aligned} \kappa_{gg} = & \frac{\alpha_3}{8\pi} \left(\sum_{i=\text{fermions}} \eta_i d_i^{(2)} C_i^{(3)} \frac{y_i}{m_i} \cdot \frac{4}{3} f_{1/2}(\tau_i) \right. \\ & \left. + \sum_{i=\text{scalars}} \eta_i d_i^{(2)} C_i^{(3)} \frac{A_i}{m_i^2} \cdot \frac{1}{6} f_0(\tau_i) \right), \end{aligned} \quad (4)$$

$$\begin{aligned} \kappa_{\gamma\gamma} = & \frac{\alpha_{\text{em}}}{8\pi} \left(\sum_{i=\text{fermions}} \eta_i \text{tr}(Q_i^2) \frac{y_i}{m_i} \cdot \frac{4}{3} f_{1/2}(\tau_i) \right. \\ & \left. + \sum_{i=\text{scalars}} \eta_i \text{tr}(Q_i^2) \frac{A_i}{m_i^2} \cdot \frac{1}{6} f_0(\tau_i) \right), \end{aligned} \quad (5)$$

in the case of scalar S , and

$$\kappa_{gg} = \frac{\alpha_3}{8\pi} \sum_{i=\text{fermions}} \eta_i d_i^{(2)} C_i^{(3)} \frac{y_{5i}}{m_i} \cdot 2\tilde{f}_{1/2}(\tau_i), \quad (6)$$

$$\kappa_{\gamma\gamma} = \frac{\alpha_{\text{em}}}{8\pi} \sum_{i=\text{fermions}} \eta_i \text{tr}(Q_i^2) \frac{y_{5i}}{m_i} \cdot 2\tilde{f}_{1/2}(\tau_i), \quad (7)$$

in the case of pseudoscalar S . Here, $d_i^{(N)}$ and $C_i^{(N)}$ are the dimension and the Dynkin index of the representation $R_i^{(N)}$ of $SU(N)$, respectively. For instance, $(d_i^{(N)}, C_i^{(N)}) = (1, 0)$, $(N, 1/2)$ and $(N^2 - 1, N)$ for $R_i^{(N)}$ being singlet, fundamental representation, and adjoint representation, respectively. The trace of the electric charge squared is given by

$$\text{tr}(Q_i^2) = d_i^{(3)} C_i^{(2)} + d_i^{(3)} d_i^{(2)} Y_i^2, \quad (8)$$

and the loop functions are (for $\tau < 1$)

$$f_{1/2}(\tau) = \frac{3}{2\tau^2} \left(\tau + (\tau - 1) \arcsin^2 \sqrt{\tau} \right), \quad (9)$$

$$\tilde{f}_{1/2}(\tau) = \frac{1}{\tau} \arcsin^2 \sqrt{\tau}, \quad (10)$$

$$f_0(\tau) = \frac{3}{\tau^2} \left(\arcsin^2 \sqrt{\tau} - \tau \right), \quad (11)$$

with $\tau_i = m_S^2/4m_i^2$. (These loop functions are normalized so that they become 1 for $\tau \rightarrow 0$.)

Let us now discuss the running gauge couplings of the SM for a scale at $\mu > m_i$, which are given by, at the one loop,

$$\alpha_a^{-1}(\mu) \simeq \alpha_{a,\text{SM}}^{-1}(m_i) - \frac{b_a^{\text{SM}} + \Delta b_a}{2\pi} \log \left(\frac{\mu}{m_i} \right), \quad (12)$$

where $\alpha_{a,\text{SM}}(m_i)$ is evaluated by using the renormalization group (RG) equations of the SM, $b_a^{\text{SM}} = (41/6, -19/6, -7)$, and

$$\begin{aligned} \Delta b_a = & \frac{4}{3} \sum_{i=\text{fermions}} \eta_i \begin{pmatrix} d_i^{(3)} d_i^{(2)} Y_i^2 \\ d_i^{(3)} C_i^{(2)} \\ d_i^{(2)} C_i^{(3)} \end{pmatrix} \\ & + \frac{1}{3} \sum_{i=\text{scalars}} \eta_i \begin{pmatrix} d_i^{(3)} d_i^{(2)} Y_i^2 \\ d_i^{(3)} C_i^{(2)} \\ d_i^{(2)} C_i^{(3)} \end{pmatrix}. \end{aligned} \quad (13)$$

Note that the contributions of each fermion or scalar to Δb_a in (13) are the same as the coefficients in the diphoton production rate, Eqs. (4)–(7). Therefore, by defining effective masses m_i^{eff} and its minimal value as

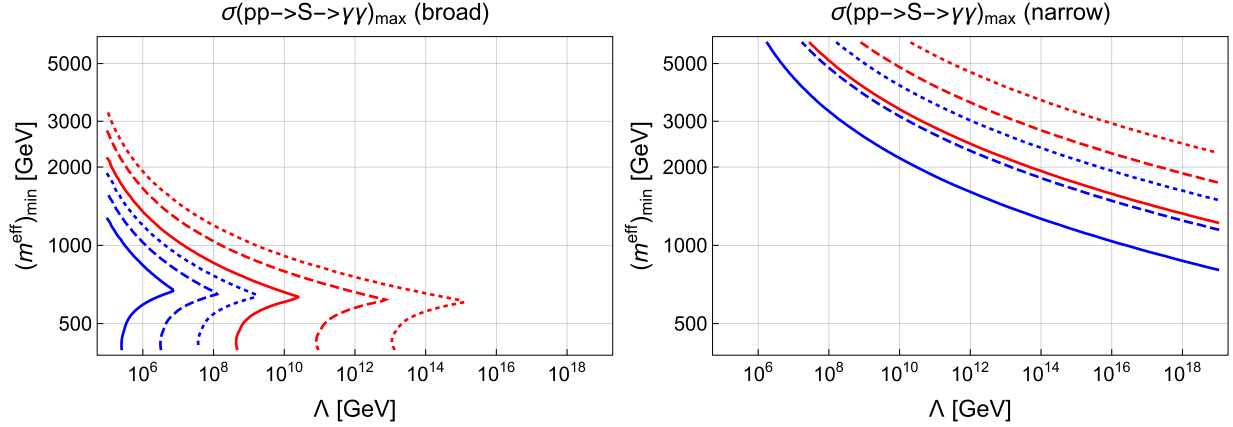


Fig. 1. Contour plots of the upper bounds on the signal rate $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)_{\max}$ as functions of the cutoff scale Λ and the minimum effective mass in the loop, m_{\min}^{eff} , defined in Eq. (14). Left: broad width case, $\Gamma_{S,\text{total}} = 45$ GeV. Right: narrow width case, $\Gamma_{S,\text{total}} = \Gamma(S \rightarrow gg) + \Gamma(S \rightarrow \gamma\gamma)$. The solid, dashed, and dotted lines show the contours of $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)_{\max} = 10, 5,$ and 3 fb, respectively. The blue lines represent the case that S is a scalar, while red lines are for the pseudoscalar case. (For interpretation of the references to color in this figure, the reader is referred to the web version of this article.)

$$m_i^{\text{eff}} \equiv \begin{cases} \frac{m_i}{y_i f_{1/2}(\tau_i)} \text{ or } \frac{m_i}{y_{5i} \tilde{f}_{1/2}(\tau_i)} & (i = \text{fermion}), \\ \frac{2m_i^2}{A_i f_0(\tau_i)} & (i = \text{scalar}), \end{cases} \quad (14)$$

$$m_{\min}^{\text{eff}} \equiv \min_i \{m_i^{\text{eff}}\},$$

one can obtain upper bounds on $\Gamma(S \rightarrow gg)$ and $\Gamma(S \rightarrow \gamma\gamma)$ as functions of Δb_a and m_{\min}^{eff} as follows;

$$\Gamma(S \rightarrow gg) < \Gamma(S \rightarrow gg)_{\max} = \frac{2m_S^3}{\pi} \left(\frac{\alpha_3}{8\pi}\right)^2 \left(\frac{\Delta b_3}{m_{\min}^{\text{eff}}}\right)^2 \times \begin{cases} 1 & (\text{scalar } S) \\ 9/4 & (\text{pseudoscalar } S) \end{cases}, \quad (15)$$

$$\Gamma(S \rightarrow \gamma\gamma) < \Gamma(S \rightarrow \gamma\gamma)_{\max} = \frac{m_S^3}{4\pi} \left(\frac{\alpha_{\text{em}}}{8\pi}\right)^2 \left(\frac{\Delta b_1 + \Delta b_2}{m_{\min}^{\text{eff}}}\right)^2 \times \begin{cases} 1 & (\text{scalar } S) \\ 9/4 & (\text{pseudoscalar } S) \end{cases}. \quad (16)$$

We consider the following two constraints on Δb_a :

- (i) **Landau pole:** We require that the SM gauge couplings are perturbative up to a scale Λ ,²

$$\alpha_a(\Lambda) < 1, \quad (17)$$

which leads to upper bounds on Δb_a as functions of Λ (and m_i).

- (ii) **Running α_3 :** In addition, too large Δb_3 (with relatively small m_i) modifies the evolution of the strong coupling constant and conflicts with the scale dependence of α_3 observed by the LHC [8]. We require that Δb_3 is below the 2σ upper bound given in Ref. [8].³ For instance, the bound is $\Delta b_3 < 5.2$ (15.9) when the mass of the particle in the loop is 500 (700) GeV.

These bounds on Δb_a lead to the maximal values of $\Gamma(S \rightarrow gg)$ and $\Gamma(S \rightarrow \gamma\gamma)$ according to Eqs. (15) and (16), which are then converted to the upper bound on the cross section for the process

$pp \rightarrow S \rightarrow \gamma\gamma$. In particular, as one can see from Eq. (1), the cross section becomes larger as $\Gamma(S \rightarrow gg)$ increases. In addition, when $\Gamma(S \rightarrow gg)$ takes its largest possible value, the partial decay rates into electroweak gauge boson pairs are always much smaller than $\Gamma(S \rightarrow gg)$, and $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ increases as $\Gamma(S \rightarrow \gamma\gamma)$ becomes larger. Thus, with Λ and m_{\min}^{eff} being fixed, the cross section takes its largest value when $\Delta b_1, \Delta b_2$ and Δb_3 are all maximized.

Fig. 1 shows the upper bound on $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ as a function of Λ and m_{\min}^{eff} , which is obtained from Eqs. (1), (15), and (16). The left figure shows the case of fixed broad width $\Gamma_{S,\text{total}} = 45$ GeV, while the right figure represents the case of narrow width, $\Gamma_{S,\text{total}} = \Gamma(S \rightarrow gg) + \Gamma(S \rightarrow \gamma\gamma)$.⁴ The red and blue lines show the cases that S is a scalar and a pseudoscalar, respectively. Here, for simplicity, we have taken $m_i = m_{\min}^{\text{eff}}$ to calculate the running coupling with Eq. (12), and also to obtain the upper bound on Δb_3 from Ref. [8].⁵

As can be seen in the left panel of Fig. 1, the cutoff scale Λ cannot be very large for a broad width case. Below the kink at $m \simeq 600\text{--}700$ GeV, the constraint from the $\alpha_3(\mu)$ measurement gives a severe upper bound on Δb_3 . In this region, the upper bound on Λ is determined by the condition of Landau poles of $\alpha_{1,2}$. Above the kink, the Landau pole condition on Δb_3 is stronger than that from the $\alpha_3(\mu)$ measurement.

In the narrow width case shown in the right panel of Fig. 1, the bounds become weaker than the broad width case, but they still constrain the region of $m_{\min}^{\text{eff}} \simeq \mathcal{O}(\text{TeV})$ when the cutoff scale Λ is large. For instance, in order to have $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 10$ fb with $\Lambda = 10^{18}$ GeV (10^{15} GeV), the effective mass should be $m_{\min}^{\text{eff}} \lesssim 870$ (1100) GeV in the case that S is a scalar, and $m_{\min}^{\text{eff}} \lesssim 1300$ (1700) GeV in the case of pseudoscalar.

Before closing this section, several comments are in order.

- The bounds in Fig. 1 are very conservative, and they can become severer in concrete and realistic models. First of all, $\Delta b_1, \Delta b_2$ and Δb_3 are simultaneously maximized in Fig. 1, but it is not generically the case in concrete models. Secondly,

² We have checked that the numerical results are almost unchanged as far as $\alpha_a(\Lambda)$ is larger than 1.

³ Although there are also similar bounds on $\Delta b_{1,2}$ from the measurements of running electroweak couplings $\alpha_{1,2}$ [10], we found that the constraints are too weak to constrain the diphoton models.

⁴ We have checked that the result does not change much even if we include other decay modes into electroweak gauge boson pairs.

⁵ As we shall see in the next section, the physical masses m_i are typically smaller than the effective mass m_i^{eff} in concrete models. Smaller masses give severer upper bounds on Δb_a for both of the constraints (i) and (ii), and hence taking $m_i = m_{\min}^{\text{eff}}$ leads to conservative constraints.

the region with small mass and large Δb_a is severely constrained by the direct search for the new particles ψ_i and ϕ_i . For instance, in the broad width case, the constraint from the $\alpha_3(\mu)$ measurement for $m = 600$ GeV is about $\Delta b_3 < 8.7$, and the upper bound corresponds to 13 Dirac pairs of vector-like quarks if they are in fundamental representations. Such a model is likely to be already excluded by direct searches, unless the new colored particles decay in a very complicated manner to escape from LHC searches. The direct search can constrain the model for the narrow width case as well. (See also the discussion in the next section.)

Although it is difficult to saturate the bounds in Fig. 1 in concrete realistic models, they are model-independent and conservative, and yet constraining interesting regions of m_{\min}^{eff} and Λ . Therefore the bounds in Fig. 1 can be an important first step to explore the physics behind the diphoton signal.

- In models with fermion loops, the Yukawa coupling $y_{(5)i}$ at low energy becomes typically smaller than unity due to the running, and hence the masses of the particles in the loop m_i should be even smaller than m_{\min}^{eff} (cf. Eq. (14)). In other words, if one adjusts the Yukawa couplings at TeV scale to larger values, the scale of the Landau pole of the Yukawa coupling becomes even smaller than those of the gauge couplings. (See the next section.)

3. Explicit examples

In the previous section, we have derived a generic upper bound on $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ for given cutoff scale Λ and effective mass scale m_{\min}^{eff} . Although it is a prominent implication of the diphoton resonance, m_{\min}^{eff} does not directly correspond to physical masses of new charged particles. In order to see how light the charged particles should be in models with fermion loops, in this section we consider the running of the Yukawa couplings with concrete examples. We also briefly discuss the LHC constraints on vector-like quarks, and comment on the running of the scalar quartic coupling of S . In this section, we only consider the case of narrow width and take $\Gamma_{S,\text{total}} = \sum_{VV=gg,\gamma\gamma,\gamma Z,ZZ,WW} \Gamma(S \rightarrow VV)$.

For simplicity, we consider the N copies of Dirac fermions which transform as $(R^{(3)}, R^{(2)}, Y)$ under the SM gauge group, with universal Yukawa coupling and mass, $y_{(5)i} = y$ and $m_i = m$. The RG equation for Yukawa coupling is given by [11]

$$16\pi^2 \frac{dy}{d \ln \mu} = \left(3 + 2d^{(2)}d^{(3)}N\right)y^3 - \left(\frac{48C^{(3)}}{d^{(3)}}g_3^2 + \frac{18C^{(2)}}{d^{(2)}}g_2^2 + 6Y^2g_1^2\right)y, \quad (18)$$

which holds both for scalar S ($y = y_i$) and pseudoscalar S ($y = y_{5i}$). For a given representation $(R^{(3)}, R^{(2)}, Y)$, one can obtain the upper bound on $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ as a function of m and Λ from the following procedure.

1. An upper bound on the multiplicity N , N_{\max} , is obtained as a function of m and Λ , by requiring that (i) the gauge couplings remain perturbative up to the scale Λ , and (ii) Δb_3 satisfies the constraint from the $\alpha_3(\mu)$ measurement [8] (see Sec. 2). For the former constraint, we require $\alpha_a(\Lambda) \leq 1$ for $a = 1-3$ in the numerical calculation.
2. For a given N ($1 \leq N \leq N_{\max}$), an upper bound on the Yukawa coupling at low energy, $y(\mu = m)$, is obtained by requiring that the running Yukawa coupling, $y(\mu)$, also remains perturbative for $\mu < \Lambda$. This gives the upper bounds on $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ for a given set of (m, Λ, N) . Because $y(m)$ increases

as $y(\Lambda)$ increases, we take $y(\Lambda) = 4$. (We have checked that the maximal possible value of the cross section does not change much as far as $y(\Lambda)$ is large enough.)

3. The maximum signal rate $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)_{\max}$ is obtained with respect to N .

In the case where there is only one representation, $N = N_{\max}$ gives the maximum value of $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ with m , Λ , and the representation of the fermion being fixed. This is because the maximal value of the Yukawa coupling at low energy roughly scales as $y \sim N^{-1/2}$, and therefore the signal rate increases as $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) \sim (Ny)^2 \sim N$.

As explicit examples, we further assume that the new vector-like fermions can decay into SM particles with a renormalizable interaction. Then, there are seven possibilities [5]

$$\begin{aligned} (R^{(3)}, R^{(2)}, Y) = & (\mathbf{3}, \mathbf{1}, -1/3), (\mathbf{3}, \mathbf{1}, 2/3), \\ & (\mathbf{3}, \mathbf{2}, 1/6), (\mathbf{3}, \mathbf{2}, -5/6), (\mathbf{3}, \mathbf{2}, 7/6), \\ & (\mathbf{3}, \mathbf{3}, -1/3), (\mathbf{3}, \mathbf{3}, 2/3). \end{aligned} \quad (19)$$

For those representations, we have numerically solved the RG equation in (18) as well as those of gauge coupling constants, and calculated the maximum signal rate $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)_{\max}$. In the following, we mainly discuss three cases $(\mathbf{3}, \mathbf{1}, 2/3)$, $(\mathbf{3}, \mathbf{2}, 7/6)$, and $(\mathbf{3}, \mathbf{3}, 2/3)$, since they have large hyper-charges and give largest signal rate among SU(2) singlets, doublets and triplets, respectively. We will briefly discuss the other cases at the end of this section.

The results for $(\mathbf{3}, \mathbf{1}, 2/3)$, $(\mathbf{3}, \mathbf{2}, 7/6)$, and $(\mathbf{3}, \mathbf{3}, 2/3)$ are shown in Fig. 2. As one can see, they look qualitatively similar to Fig. 1. However, with m_{\min}^{eff} being fixed, the cross section is smaller than that given in Fig. 1 because Δb_1 and Δb_2 are not simultaneously maximized. Furthermore, m is smaller than m_{\min}^{eff} because the Yukawa coupling becomes smaller than unity at low energy, in particular when N is large (cf. Eq. (18)).

- Among the three cases, $(\mathbf{3}, \mathbf{2}, 7/6)$ gives the largest $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)_{\max}$ in most of the parameter space. For instance, in the pseudoscalar case, $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 10$ fb can be realized with a cutoff scale of 10^{16} GeV (10^{10} GeV) if the fermions are lighter than 740 GeV (1100 GeV). For $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 5$ fb, the fermions can be as heavy as 1000 GeV (1600 GeV) for the cutoff scale of 10^{16} GeV (10^{10} GeV).
- In the case of SU(2) triplet $(\mathbf{3}, \mathbf{3}, 2/3)$, the maximal signal rate $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)_{\max}$ is smaller than the case of $(\mathbf{3}, \mathbf{2}, 7/6)$ in most of the parameter space, except for the large cutoff region $\Lambda \gtrsim 10^{17}$ GeV. In this case, the cutoff above the Planck scale is allowed, e.g., for the pseudoscalar case with $N = 1$, $m \simeq 860$ GeV, and $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 5$ fb.
- Finally, in the case of SU(2) singlet $(\mathbf{3}, \mathbf{1}, 2/3)$, the signal rate is suppressed compared with the other two cases. In this case, the running α_3 constraint determines the N_{\max} in a large part of the low mass region $m \lesssim 600-700$ GeV. The zigzag lines for $m \lesssim 600$ GeV is due to the rapid increase of allowed N_{\max} (Δb_3) with respect to m from the running α_3 constraint. In each narrow range of m with a fixed N_{\max} , the upper bound on Λ is determined either by the perturbativity of the Yukawa coupling or by the Landau pole of U(1) $_Y$.

Next, we discuss the constraint from the direct searches for vector-like quarks at the LHC. Here, we assume that they decay into the SM particles via a renormalizable coupling with SM quarks and the Higgs boson. In order to avoid the stringent constraint from the decay into third generation quarks, let us further assume that the coupling with the third generation is suppressed.

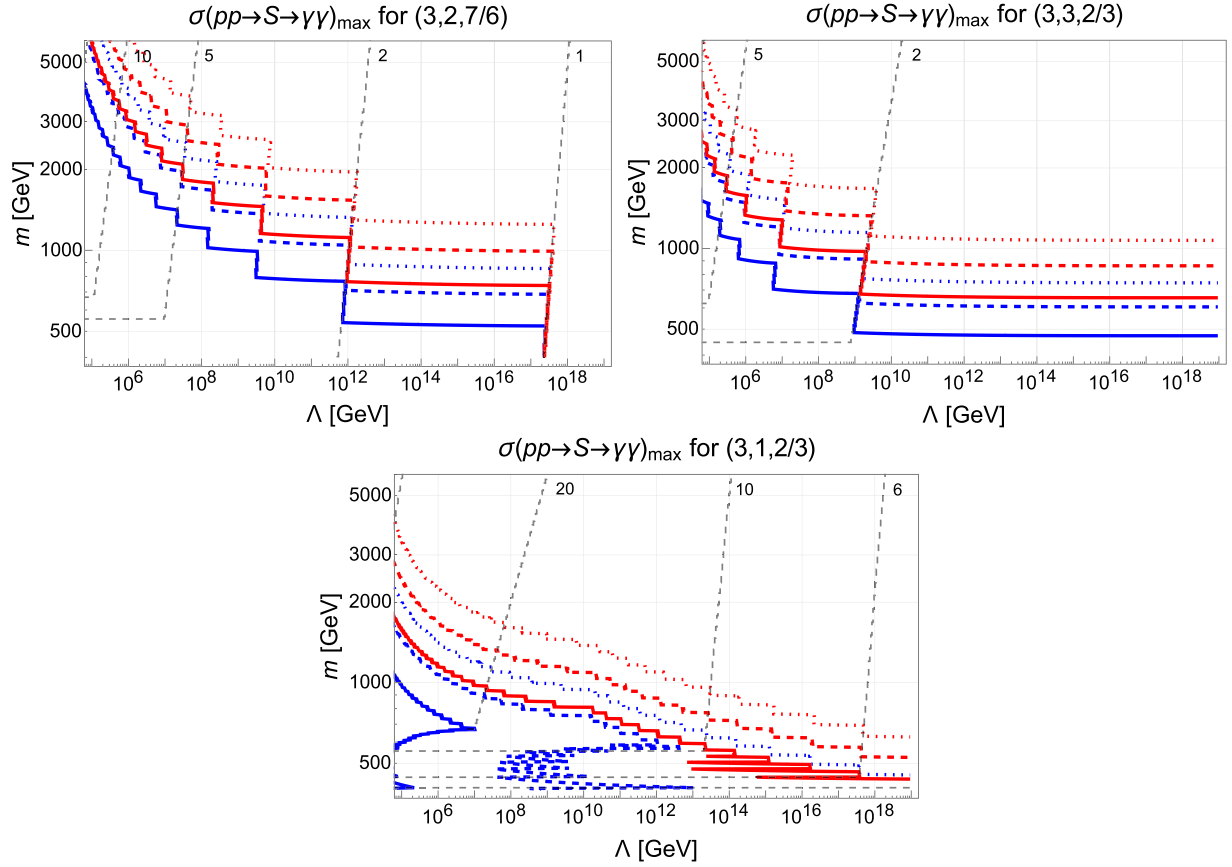


Fig. 2. The upper bound on the signal rate $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)_{\max}$ as a function of the cutoff scale Λ and the fermion mass m in the case of Dirac fermions with quantum numbers of $(\mathbf{3}, \mathbf{2}, 7/6)$ (top left), $(\mathbf{3}, \mathbf{1}, 2/3)$ (bottom), and $(\mathbf{3}, \mathbf{3}, 2/3)$ (top right). Here, we take $\Gamma_{S,\text{total}} = \sum_{V=gg,\gamma\gamma,\gamma Z,ZZ,WW} \Gamma(S \rightarrow VV)$. The solid, dashed, and dotted lines show the contours of $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)_{\max} = 10, 5,$ and 3 fb, respectively. The blue lines represent the case that S is a scalar, while red lines are for the pseudoscalar case. Black dashed lines show the maximal allowed number of generations, N_{\max} . (For interpretation of the references to color in this figure legend, the reader is referred to the web version of this article.)

Then, the vector-like quarks decay into a light SM quark and a W/Z /Higgs boson, depending on its representation. In particular, the search for a vector-like quark decaying into a W boson and a light SM quark at the LHC gives a stringent constraint in the present scenario. From the result of ATLAS [12], the bound is estimated as⁶

$$\begin{cases} N \cdot \text{Br}(Q' \rightarrow Wq) \lesssim 1 & \text{for } m \leq 690 \text{ GeV}, \\ N \cdot \text{Br}(Q' \rightarrow Wq) \lesssim 1.5 & \text{for } 690 \text{ GeV} \leq m \lesssim 750 \text{ GeV}, \\ N \cdot \text{Br}(Q' \rightarrow Wq) \lesssim 2 & \text{for } 750 \text{ GeV} \lesssim m \lesssim 800 \text{ GeV}, \end{cases} \quad (20)$$

where Q' and q denote the vector-like quark and the SM light quark, respectively.

- In the case of the SU(2) doublet $(\mathbf{3}, \mathbf{2}, 7/6)$, it contains vector-like quarks with electric charges of $5/3$ and $2/3$. The one with the electric charge of $5/3$ decays into a W boson and a light SM quark (up and/or charm) with almost 100% branching fraction. Comparing the bound in (20) with the lines in Fig. 2, if we require $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 10$ fb in the case of scalar S , the region of $\Lambda \gtrsim 10^{12}$ GeV ($N = 1$) is excluded, and $\Lambda \gtrsim 10^{9.5}$ GeV ($N = 2$) is at the boundary of excluded region. In the case of pseudoscalar, the model can explain $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 10$ fb while being perturbative up to

$\Lambda \simeq 10^{17}$ GeV ($\Lambda \simeq 10^{12}$ GeV), if the vector-like quarks are as light as about 740 GeV (1100 GeV).

- In the case of SU(2) triplet $(\mathbf{3}, \mathbf{3}, 2/3)$, one of the SU(2) triplet quarks decays into Wq with an almost 100% branching fraction, and another one has about 50% branching. Thus, from the bound (20), the region of $m \lesssim 750$ GeV is excluded even for $N = 1$. The scalar case cannot have a cutoff larger than about 10^9 GeV in order to have the cross section larger than ~ 3 fb, while the pseudoscalar case with $N = 1$, $m \simeq 860$ GeV, and $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 5$ fb is still allowed and can be perturbative up to the Planck scale.
- Finally, in the case of SU(2) singlet $(\mathbf{3}, \mathbf{1}, 2/3)$, the direct search excludes a large fraction of the parameter space with a sizable signal cross section, in particular when the cutoff scale is high. In this case, the vector-like quark decays into a W boson and a light quark with a branching fraction of about 50%. From the direct search bound in Eq. (20), the number of multiplicity should satisfy $N < 2, 3,$ and 4 for $m \leq 690$ GeV, $m \simeq (690-750)$ GeV, and $m \simeq (750-800)$ GeV, respectively. Thus, the lines in the figure for $m \lesssim 800$ GeV are not consistent with the direct search bound. If we adopt the maximal number of multiplicity allowed by the direct search, the Yukawa coupling should be quite large at low energy in order to explain the diphoton signal. Even for the pseudoscalar case and for $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 3$ fb, the required value of the Yukawa coupling is $y(m) \gtrsim 2.1, 1.5,$ and $1.2,$ for $m \leq 690$ GeV, $m \simeq (690-750)$ GeV, and $m \simeq (750-800)$ GeV, respectively. If the RG equation (18) is evolved from low energy to high

⁶ We could not find any constraint for vector-like quarks with $m > 800$ GeV decaying into light quarks. Thus, we do not consider the direct bound for $m > 800$ GeV.

energy, they quickly become non-perturbative, which leads to cutoff scales below 10 TeV.

We should note that the above constraints strongly depend on the decay modes of vector-like quarks. If they mainly couple to the third generation SM quarks and decay into top and/or bottom quarks, the constraints become severer. Instead, if they decay in a very complicated way (e.g., in a cascade decay chain with multiple intermediate new particles emitting many soft jets), they may escape the direct search even for small mass region.

Now let us briefly discuss the other representations in Eq. (19).

- In the case of (3, 1, $-1/3$), we checked that, even for the pseudoscalar case, $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ is smaller than 3 fb for $\Lambda > 10^5$ GeV.
- In the case of (3, 2, $1/6$), the pseudoscalar case can have $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = (5\text{--}10)$ fb with a large cutoff, but it requires a small mass m and a large multiplicity N . We found that the region below $m \lesssim 800$ GeV is excluded if the vector-like quarks mainly decay into the SM light quarks, and for $m \gtrsim 800$ GeV the cutoff cannot be larger than 10^9 GeV for $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) \geq 5$ fb.
- The case of (3, 2, $-5/6$) is similar to that of (3, 2, $7/6$), but with the multiplicity N roughly twice as large. For instance, $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) \simeq 5$ fb can be obtained in the pseudoscalar case with $N = 2$, $m \simeq 910$ GeV, and the cutoff scale as large as $\Lambda \simeq 10^{17}$ GeV.
- The case of (3, 3, $-1/3$) is similar to (3, 3, $2/3$), but with the fermion mass m being about 30% smaller.

Before closing this section, we comment on the running of the quartic coupling of the S field.⁷ Defining the coupling λ as $\mathcal{L}_{S^4} = -(1/4)\lambda S^4$, its RG equation is given by

$$16\pi^2 \frac{d\lambda}{d \ln \mu} = 3\lambda^2 - 48Nd^{(3)}d^{(2)}y^4 + 8Nd^{(3)}d^{(2)}y^2\lambda. \quad (21)$$

We have checked that, as far as λ is positive at the cutoff scale, it does not become negative for $m < \mu < \Lambda$ and hence there is no vacuum instability. In addition, λ does not blow up below the cutoff scale irrespective of the value $\lambda(\Lambda)$. Thus, there is no constraint from the running of the quartic coupling.

4. Conclusions

Motivated by the recent LHC results, we have studied the diphoton resonance production cross section at the LHC, paying particular attention to the running of the gauge and Yukawa coupling constants. We have considered the case where a (pseudo-) scalar particle S with its mass of 750 GeV is responsible for the diphoton events observed by the LHC and the scalar particle is produced by the gluon fusion. In such a case, new fermions and/or bosons which have SM gauge quantum numbers are necessary to generate S - g - g and S - γ - γ vertices. Assuming that the S - g - g and S - γ - γ vertices are perturbatively generated by the loop effects of the new fermions and/or bosons, we studied how large the cross section for the process $pp \rightarrow S \rightarrow \gamma\gamma$ can be. We have shown that the cross section is severely constrained from above by (i) the perturbativity of the coupling constants up to a certain scale, and (ii) the consistency of the scale dependence of α_3 with that observed by the LHC.

First, we have pointed out that a model-independent upper bound on $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ can be derived, taking account of

the two requirements mentioned above. Such a bound is obtained from the fact that the cross section is related to $\Gamma(S \rightarrow gg)$ and $\Gamma(S \rightarrow \gamma\gamma)$, and that the amplitudes for these decay rates are proportional to the β -function coefficients of the gauge coupling constants from the fermions and bosons inside the loop. We have also calculated such a bound as a function of the cutoff scale Λ and the m_{\min}^{eff} parameter which corresponds to the mass scale of the fermions and bosons inside the loop. (See Fig. 1.)

Then, we have discussed the upper bound on $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ in models with fermion loops, taking into account the perturbativity of the Yukawa coupling between S and the new fermions. For such a study, the particle content should be fixed to perform the RG analysis. We have considered seven possible representations of the fermions with which the fermions can directly decay into SM particles. We have introduced N copies of fermions in the same representation with the universal mass of m , and derived the upper bounds on $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$. Among them, the representation of (3, 2, $7/6$) can give the largest diphoton rate in most of the parameter region. For instance, in the case of pseudoscalar, it is shown that $\sigma(pp \rightarrow S \rightarrow \gamma\gamma) = 5$ and 10 fb can be obtained with $m \simeq 1000$ and 740 GeV ($m \simeq 1600$ and 1100 GeV) and $N = 1$ ($N = 2$) when the cutoff scale is 10^{16} GeV (10^{10} GeV), respectively. We have also discussed that such sets of parameters are consistent with the current constraints on vector-like quarks from the direct search at the LHC. In the cases of the other representations, the signal rate $\sigma(pp \rightarrow S \rightarrow \gamma\gamma)$ is more suppressed, and a large cutoff scale is impossible at all in some cases.

The present study suggests that, unless the cutoff scale is very low, there must exist new particles at TeV scale or lower. They should be an important target of the LHC run-2 and other future collider experiments.

Acknowledgements

The authors thank Masafumi Kurachi for bringing our attention to the running α_3 constraint. This work was supported by Grant-in-Aid for Scientific research Nos. 23104008 (TM), 25105011 (ME), 26104001 (KH), 26104009 (KJB and KH), 26247038 (KH), 26400239 (TM), 26800123 (KH), and by World Premier International Research Center Initiative (WPI Initiative), MEXT, Japan.

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