PHYSICAL REVIEW D **94,** 055009 (2016)

Perturbativity, vacuum stability, and inflation in the light of 750 GeV diphoton excess

Mansi Dhuria*

Theoretical Physics Division, Physical Research Laboratory, Navrangpura, Ahmedabad 380 009, India

Gaurav Goswami†

Institute of Engineering and Technology, Ahmedabad University, Navrangpura, Ahmedabad 380 009, India

(Received 1 March 2016; published 9 September 2016)

The recent observation of the 750 GeV diphoton excess at 13 TeV LHC has motivated many scenarios of physics beyond the standard model. In this work, we begin by showing that many models which explain the observed excess tend to get strongly coupled well below the Planck scale. We then study a simple scenario involving colored vectorlike fermions with exotic charges, which is expected to stay weakly coupled till the Planck scale. We find the conditions under which this happens, derive the renormalization group equations for such models and solve them to show that perturbativity till Planck scale can be maintained for a very reasonable choice of parameters. Finally, we discuss issues related to vacuum stability and the possibility of inflation in the scenarios we study.

DOI: 10.1103/PhysRevD.94.055009

I. INTRODUCTION

Though the existence of physics beyond the standard model (SM) of elementary particle physics is a widespread belief, new physics is yet to show up at accessible energies at the colliders. However, recently, both the CMS and ATLAS collaboration have recently reported an excess in the diphoton channel at around 750 GeV from the data collected at the LHC run 2 at energy $\sqrt{s} = 13$ TeV [1,2]. The ATLAS has analysed 3.2 fb^{-1} of the data collected at 13 TeV, reporting an excess of about $14\gamma\gamma$ events peaked at $M = 747$ GeV. With this, the best fit value of cross-section suggests high width $\Gamma/M \sim 0.06$ with a local significance of 3.9σ. Similarly, the CMS experiment has analyzed 2.6 fb^{-1} of the data collected at 13 TeV and reported an excess of about $10\gamma\gamma$ events peaked at $M = 760$ GeV. The best fit value of the cross-section in this case suggests a very narrow width with a local significance of 2.6σ. However, the CMS best-fit width matches with the value provided by ATLAS by considering the local significance of the excess at CMS up to 2.0σ. The CMS has also presented the compatibility of the excess at both 8 TeV [3] and 13 TeV [2], suggesting a production cross section times branching ratio into photons to be 4.47 ± 1.86 fb at CMS [2]. In case of ATLAS experiment, the compatibility of excess at both 8 TeV [4] and 13 TeV [1] suggests production cross section times branching ratio into photons to be 10.6 ± 2.9 fb [5].

Although the observation requires further analysis and much more data to ensure that the excess signal is not just a statistical fluctuation, it is still tempting to study phenomenological implications of this result [5–96]. The observed resonance implies the existence of a spin-0 or spin-2 state since a spin-1 state cannot decay into two photons due to the Landau-Yang theorem [97,98].

We assume that the observed excess is due to SM singlet spin-zero particle. Since the singlet scalar does not have any direct interaction with photons, it has to decay into two photons through a loop of electrically charged particles. There exists a vast literature on diphoton excess in which the particles running in the loop are assumed to be vectorlike fermions [10,12,15]. Thus, 750 GeV diphoton excess motivates a beyond standard model (BSM) scenario involving the following particles: all SM particles, a SM singlet scalar and charged vectorlike fermions. Moreover, the existence of vectorlike fermions can in general be motivated by many theories of BSM physics, such as string-theoretic models [99,100], little Higgs models [101], and composite Higgs models [102], etc.

It is noteworthy that the search for the vectorlike fermions at LHC run 2 has pushed the limit of vectorlike quarks up to around TeV [103–105]. Motivated by this, we consider the mass of these new fermions to be 1.2 TeV. We show that the compatibility of the observed diphoton signal rate generically gives a very high value of Yukawa coupling of real singlet scalar with vectorlike fermion, thus making the theory strongly coupled. We then show that the issue can be resolved either by considering additional copies of vectorlike fermions or/and by choosing colored vectorlike fermions to have exotic charges (i.e. $Q > 2/3$).

It is well known that if SM is assumed to hold good till arbitrarily high energies, the Higgs potential turns negative at high field values (above 10^{10} GeV) suggesting the

*mansi@prl.res.in
†gaurav.goswami@ahduni.edu.in

existence of new physics beyond the SM. It is then very interesting to study the impact of the new particles needed to explain the diphoton excess on the electroweak vacuum stability of the SM. We present a detailed study of the effect of these extra particles on the scalar potential of the theory. E.g. we find the conditions under which all couplings can be kept small and positive till Planck scale. We then look into the possibility of the new scalar being responsible for inflation, similar to the SM Higgs. In order to study the consequences of this model at high energies in detail, we have found the renormalization group evolution equations for scenarios involving extra vectorlike fermions with exotic charges.

The outline of the article is as follows: In Sec. II, we begin by finding the generic conditions under which the theory remains perturbative and weakly coupled if the diphoton excess is explained by the decay of a SM singlet scalar by considering heavy colored vectorlike fermion(s) in the loop. In Sec. III, we present two simple models with exotic charge fermions which could not only explain the diphoton excess but also maintain perturbativity up to Planck scale. To achieve this, we ensure, among other things, that the mixing of the newly added scalar with SM Higgs is small and that none of the couplings, such as the gauge coupling g_1, become too large below Planck scale. In Sec. IV, we show that the proposed models (a) stay weakly coupled till Planck scale, (b) do not suffer from the vacuum stability problem, and, (c) are such that both SM Higgs and the 750 GeV scalar could act as the inflaton. For this, we solve the renormalization group (RG) evolution equations for all couplings involved in the two models. Moreover, we find that only one of the two models is consistent with the requirement of not leaving any dangerous relics in the Universe. Finally, in Sec. V, we summarize our results. In Appendix A we provide RG equations generalized to the scenarios which we study in this paper.

II. THE VARIOUS CONSTRAINTS

We interpret the diphoton signal as being due to a 750 GeV scalar (as opposed to a pseudoscalar) S produced by gluon fusion and then decaying into two photons. The cross section for this process is given by [43],

$$\sigma(gg \to S \to \gamma\gamma) = \frac{\pi^2}{8m_s^3}\mathcal{I}_{\rm pdf}\Gamma(S \to \gamma\gamma), \tag{1}$$

where, m_s is the mass of the scalar and $\mathcal{I}_{\rm pdf}$ is the dimensionless parton distribution function integral evaluated at $\sqrt{s} = 13$ TeV. Let $g_{s\gamma}$ be the coefficient of the dimension five operator $\phi F^{\mu\nu}F_{\mu\nu}$ in the low energy effective field theory. Then, the decay width $\Gamma(S \to \gamma\gamma)$ is given by $\Gamma(S \to \gamma\gamma) = g_{s\gamma}^2 m_s^3/(8\pi)$. Using this and $\mathcal{I}_{\rm pdf} \approx 5.8$ (by assuming $m_S = 750$ GeV, see [43]), we can find $g_{s\gamma}$ in terms of the total cross section $\sigma(pp \to S \to \gamma\gamma)$ as

$$\frac{g_{s\gamma}}{(\mathrm{TeV})^{-1}} = 3 \times 10^{-3}\left(\frac{\sigma}{\mathrm{fb}}\right)^{1/2}. \tag{2}$$

The cross section $\sigma(pp \to S \to \gamma\gamma)$ is inferred by ATLAS and CMS experiments to be in the range $[3 - 13]$ fb at $\sqrt{s} = 13$ TeV. Using this and Eq (2), the inferred value of the dimensionful coupling $g_{s\gamma}$ lies in the range $(5.2 \times 10^{-3},\ 1.1 \times 10^{-2})$.

At the microscopic level, the interaction of the scalar with photons is expected to be induced by loop involving other particles. In this work, we assume those particles to be fermions. The effective dimensionful coupling $g_{s\gamma}$ can then be evaluated from first principles and for a loop involving Dirac fermions in fundamental representation of $SU(3)$, coupled to the scalar S with Yukawa coupling y_M and having electric charges Q_f and masses m_f, is given by [106]

$$g_{s\gamma} = \sum_f \frac{\alpha}{4\pi}\left(\frac{2N_c y_M Q_f^2}{m_f}\right)A_{1/2}(T_f), \tag{3}$$

where, the summation is over all such fermions, α is the fine structure constant, $N_c = 3$, $T_f = 4m_f^2/m_s^2$, $A_{1/2}(x) = 2x[1 + (1 - x)f(x)]$ and $f(x)$ is of the form

$$f(x) = \begin{cases} \left(\sin^{-1}\left(\frac{1}{\sqrt{x}}\right)\right)^2, & x \geq 1, \\ -\frac{1}{4}\left(\ln\left(\frac{1+\sqrt{1-x}}{1-\sqrt{1-x}}\right) - i\pi\right)^2, & x < 1. \end{cases}$$

We are considering the scalar to be SM singlet and fermions to be vectorlike. The CMS collaboration has analyzed the data collected at LHC at $\sqrt{s} = 8$ TeV for the search of a pair production of charge $2/3$ vectorlike fermions and reported the lower bound on the mass of the same to be around 0.85 TeV [104]. On the other hand, the analysis performed by the same collaboration using the data collected at $\sqrt{s} = 13$ TeV for the search of pair-produced $5/3$ charged top partners gave a lower bound on the mass of the same to be 0.95 TeV [103]. Also, recently the ATLAS collaboration has analyzed the data collected at LHC at $\sqrt{s} = 8$ TeV for the search of singly produced charge $2/3$ or (-4/3) vectorlike fermions and provided the lower bound on the mass of the same to be around 0.95 TeV [105]. In view of this, we consider the mass of vectorlike fermion to be 1.2 TeV in our analysis. It is easy to see that, in Eq. (3) if one assumes $m_f = 1.2$ TeV, one gets $T_f = 10.24$ and $A_{1/2}(T_f) = 1.365$. Thus, we find that

$$\frac{g_{s\gamma}}{(\mathrm{TeV})^{-1}} = (3.963 \times 10^{-3})y_M Q^2, \tag{4}$$

where $Q^2 = \sum Q_f^2$. Comparison with Eq. (2) tells us that the range of values of $y_M Q^2$ which is compatible with observations lies in the range $(1.32,\ 2.73)$.

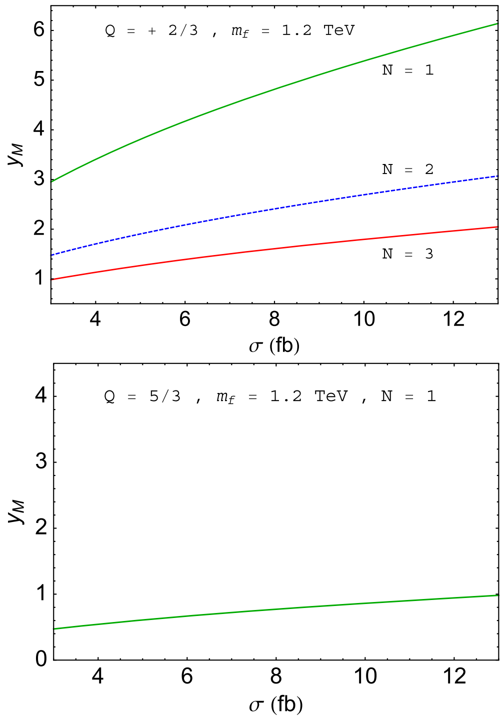

FIG. 1. The calculated values of Yukawa coupling y_M as a function of the cross section $\sigma(pp \to \gamma\gamma)$ for vectorlike fermions with up-type quark charges (top) and for vectorlike fermions with exotic charges (bottom). Notice the difference in range of y_M values in the two plots. It is easy to see that exotic charges easily restore perturbativity while increasing the number of fermions has a more mild effect on perturbativity.

In Fig. 1, we have shown the range of values of Yukawa coupling y_M which give cross section $\sigma(pp \to \gamma\gamma) \sim [3-13]$ fb by considering vectorlike fermions (with up-type quark charges) in the loop with $m_f = 1.2$ TeV. It turns out that by considering single vector fermion of electric charge $Q = 2/3$, the theory is strongly coupled for $\sigma(pp \to \gamma\gamma) \geq 4.5$ fb at the scale m_f itself. If we have N identical fermions with $Q = 2/3$ in Fig. 1, one might think that one could restore the perturbativity for $\sigma(pp \to \gamma\gamma) \sim [3-13]$ fb by considering $N = 3$. As can be checked explicitly, though this renders the Yukawa coupling sufficiently small at TeV scale, the theory still becomes nonperturbative as we approach higher energies. We will come back to this in Sec. IV. Thus it is a nontrivial task to explain the diphoton excess without running into nonperturbative regime at the TeV scale or some sub-Planckian scale.

An alternative option is to consider fermions running in the loop to have larger values of the electric charge [12]. It is easy to see from the plot at the bottom in Fig. 1 that for $Q = 5/3$, the corresponding Yukawa coupling remains small for the entire range of $\sigma(pp \to \gamma\gamma) \sim [3-13]$ fb and $N = 1$ at least at the TeV scale.

Before we proceed, let us list the various constraints which we would like to satisfy:

(1) The cross section $\sigma(pp \to S \to \gamma\gamma)$ must lie in the observed range of 3–13 fb,
(2) Perturbativity of all the couplings e.g. Yukawa, scalar couplings, gauge couplings: All couplings must be sufficiently small so that the theory stays perturbative till Planck scale,
(3) Vacuum stability: At all scales the scalar couplings must be such that the scalar potential V should always be above its value for the standard model vacuum,
(4) Mixing angle between Higgs and the singlet scalar must be such that no observable effect takes place,
(5) Given that there is going to be another scalar in the theory apart from SM Higgs, we would like to check whether one can get successful inflation in such scenarios.

In the next section, we discuss how by considering vectorlike fermions with an exotic charge $Q > 2/3$, the theory can be kept perturbative even up to Planck scale for a reasonable choice of parameters.

III. A SPECIFIC SCENARIO

The gauge group of our model is assumed to be that of the standard model (SM) i.e. $SU(3)_c \times SU(2)_L \times U(1)_Y$. The bosonic fields in our model are: the SM gauge fields, the SM Higgs field and a real parity even scalar field χ (assumed to be a singlet). We consider two models which we shall call model X and model Y and which differ in their fermionic content. In the fermionic sector of model X, there is, in addition to the SM fermions, a single colored vectorlike fermion, Ψ, with exotic electric charge which is a singlet under $SU(2)_L$. On the other hand, in model Y, apart from the SM fermions, there is colored vectorlike fermion which is a doublet under $SU(2)_L$. The reason for choosing $SU(2)_L$ doublet fermion is that it allows mixing of exotic charged vectorlike fermions with SM quarks, therefore causing decay of the same into SM particles (see Sec. IV C for details).

A. Constraints from the scalar sector

The bosonic sectors of the two models are the same. In particular, in the scalar sector, the real scalar field χ does not couple to any of the gauge fields while the Higgs field couples to the $SU(2)_L$ and $U(1)_Y$ gauge fields. The scalar potential is given by

$$V(\Phi,\chi) = \lambda_H \left(\Phi^\dagger\Phi - \frac{v^2}{2}\right)^2 + \lambda_S(\chi^2 - u^2)^2 + \lambda_{HS}\left(\Phi^\dagger\Phi - \frac{v^2}{2}\right)(\chi^2 - u^2), \tag{5}$$

so that the potential is specified by five parameters: $u, v, \lambda_H, \lambda_S, \lambda_{HS}$. Expanding the fields about the vev

$$\Phi = \frac{1}{\sqrt{2}}(0 \quad v+h)^T, \qquad \chi = u + s, \tag{6}$$

the terms quadratic in the fields h and s are

$$V_2 = (\lambda_H v^2)h^2 + (\lambda_S u^2)s^2 + (\lambda_{HS} uv)sh, \tag{7}$$

so that the mass eigenstates H and S will be different from the fields h and s and there will be a mixing between the two. The mixing angle θ and the mass eigenvalues m_H and m_S can be readily determined from the parameters specifying the potential (see e.g. [107] for details). The mixing between Higgs and new singlet S is given by [107]

$$\theta = \frac{1}{2}\tan^{-1}\left(\frac{\lambda_{SH} u\, v}{\lambda_S u^2 - \lambda_H v^2}\right). \tag{8}$$

To avoid the tight constraints from the LHC on the Higgs decay modes, it a good idea to suppress the mixing between S and H. We thus need to ensure that the parameters λ_H, λ_S, λ_{HS} are such that for $u \sim \mathcal{O}$ (TeV), θ turns out to be less than 0.15 [108] and m_H, m_S take up the expected values.

As our scenario involves a heavy scalar, at the scale m_S, we can integrate out this heavy scalar and hence solve the equations of motion of the theory to obtain a description of the potential in terms of the field Φ alone. This causes a discontinuity in λ_H at the scale m_S so that the new value of the quartic coupling is identified with $\lambda_H^{\rm SM}$, the Higgs self-coupling in SM [109] i.e. at the scale m_S,

$$\lambda_H - \frac{\lambda_{SH}^2}{4\lambda_S} = \lambda_H^{\rm SM}. \tag{9}$$

B. Constraints from fermionic sector

In both models, the newly added fermions are colored so that they couple to the gluons. In addition to this, in model X, the fermions couple to the $U(1)_Y$ gauge field while in model Y, the fermions couple to the gauge fields of both $SU(2)_L$ and $U(1)_Y$. This ensures that in model Y, we will have terms such as $g_2\bar{\psi}_2\gamma^\mu W_\mu^- \psi_1$ and $g_2\bar{\psi}_1\gamma^\mu W_\mu^+ \psi_2$. The presence of these terms plays a crucial role in this model.

Now let us turn to the Yukawa terms in the two models. In what follows, the quarks always refer to the third generation quarks of the SM. In model X, the Yukawa term of the form $\chi\bar{\psi}_L\psi_R$ is possible irrespective of the charge of the fermion. Moreover, the terms $\chi\bar{\psi}_L u_R$ and $\Phi^c\bar{Q}_L\psi_R$ are also possible if the charge of the fermion is chosen to be $2/3$. Similarly, the terms $\chi\bar{\psi}_L d_R$ and $\Phi\bar{Q}_L\psi_R$ are also possible if the charge of the fermion is chosen to be $-1/3$.

In model Y, we could denote the doublet fermion ψ by $\psi = (\psi_1\psi_2)^T$. Since the hypercharge of all the fields which are part of the same multiplet is the same, it is easy to see that $Q_1 - Q_2 = 1$. In model Y, the Yukawa term of the form $\chi\bar{\psi}_L\psi_R$ is possible irrespective of the hypercharge of the fermion. Moreover, the terms $\Phi\bar{\psi}_L d_R$, $\Phi^c\bar{\psi}_L u_R$ and $\chi\bar{Q}_L\psi_R$ are also possible if the hypercharge of the doublet fermion is chosen to be $1/6$. In this case the electric charges of the particles in the doublet are $Q_1 = 2/3$ and $Q_2 = -1/3$, the same as that in the SM quark doublet. Similarly, the term $\Phi\bar{\psi}_L u_R$ is possible if the hypercharge of the doublet is $7/6$, this gives $Q_1 = 5/3$ and $Q_2 = 2/3$. Finally, the term $\Phi^c\bar{\psi}_L d_R$ is possible if the hypercharge of the doublet is $-5/6$, which leads to $Q_1 = -1/3$ and $Q_2 = -4/3$.

The list of all possible Yukawa couplings with different hypercharges in both the models is summarized in Table I. It is easy to verify that this exhausts all the possibilities for BSM Yukawa terms in both model X and model Y. The arguments of Sec. II suggest that the charge of the newly added fermion must be exotic. This means that the Yukawa part of the Lagrangian in model X is

$$\mathcal{L}_{\rm Yuk}^{\rm X} = \mathcal{L}_{\rm Yuk}^{\rm SM} + y_M\chi\bar{\psi}_L\psi_R + {\rm H.c.}, \tag{10}$$

while the Yukawa part of the Lagrangian for the case with $7/6$ hypercharge in model Y is

$$\mathcal{L}_{\rm Yuk}^{\rm Y} = \mathcal{L}_{\rm Yuk}^{\rm SM} + y_M\chi\bar{\psi}_L\psi_R + y_f\Phi\bar{\psi}_L u_R + {\rm H.c.}, \tag{11}$$

and the Yukawa part of the Lagrangian for the case with $-5/6$ hypercharge in model Y is

$$\mathcal{L}_{\rm Yuk}^{\rm Y} = \mathcal{L}_{\rm Yuk}^{\rm SM} + y_M\chi\bar{\psi}_L\psi_R + y_f'\Phi^c\bar{\psi}_L d_R + {\rm H.c.}. \tag{12}$$

TABLE I. Possible Yukawa terms in the Lagrangians of model X and model Y and the corresponding hypercharge of the newly added vectorlike fermion.

Model	Yukawa terms	Hypercharge Y_f	Vector-like fermion content
X	$\chi\bar{\psi}_L\psi_R$ $\chi\bar{\psi}_L u_R$ $\Phi^c\bar{Q}_L\psi_R$	2/3	$\psi_{2/3}$
	$\chi\bar{\psi}_L\psi_R$ $\chi\bar{\psi}_L d_R$ $\Phi\bar{Q}_L\psi_R$	-1/3	$\psi_{-1/3}$
	$\chi\bar{\psi}_L\psi_R$	Exotic/Arbitrary ($Y_f > 2/3$)	$\psi_{(Q>2/3)}$
Y	$\chi\bar{\psi}_L\psi_R$ $\Phi\bar{\psi}_L d_R$ $\Phi^c\bar{\psi}_L u_R$ $\chi\bar{Q}_L\psi_R$	1/6	$(\psi_{2/3}\psi_{-1/3})^T$
	$\chi\bar{\psi}_L\psi_R$ $\Phi\bar{\psi}_L u_R$	7/6	$(\psi_{5/3}\psi_{2/3})^T$
	$\chi\bar{\psi}_L\psi_R$ $\Phi^c\bar{\psi}_L d_R$	-5/6	$(\psi_{-1/3}\psi_{-4/3})^T$

C. Issues of perturbativity: Gauge couplings

By now it is clear that we shall be interested in scenarios in which vectorlike fermions with exotic charges are added to the SM. Each gauge coupling g_i evolves as (see Appendix)

$$\frac{dg_i}{dt} = \frac{b_i}{(4\pi)^2} g_i^3, \tag{13}$$

where b_i are numerical factors. Just like in the SM, in both model X and model Y, both b_3 and b_2 are negative, so, g_3 and g_2 become small at large energies. In the models we are considering

$$b_1 = \begin{cases} \frac{41}{10} + \frac{12}{5} Y^2, & \text{model X}, \\ \frac{41}{10} + \frac{24}{5} Y^2, & \text{model Y}, \end{cases} \tag{14}$$

is a positive quantity. The value of b_1 is larger in both the models X and Y as compared to the SM in which $b_1 = 4.1$. If the coupling g_1 at the scale μ_i is $g_1(\mu_i)$, then g_1 at a higher scale μ is

$$g_1(\mu) = \frac{g_1(\mu_i)}{(1 - \frac{b_1 g_1(\mu_i)^2}{8\pi^2} \ln(\frac{\mu}{\mu_i}))^{1/2}}, \tag{15}$$

so that, at the scale

$$\mu = \mu_i \exp\left(\frac{8\pi^2}{b_1 g_1(\mu_i)^2}\right), \tag{16}$$

the coupling g_1 hits a Landau pole. For SM, this happens at $\mu \approx 10^{41}$ GeV (in one-loop approximation). In the scenarios we are interested in, larger values of b_1 turn up and the Landau pole is found at a lower scale.

Now, one may wish to have inflation with the scalars present in the model. The potential at large values of the fields is quartic and this leads to large field inflation. During large field inflation, the inflaton field undergoes super-Planckian excursion so that to get the minimum 60 e-foldings of inflation, the field excursion $\Delta\phi \sim \mathcal{O}(10) M_{\rm Pl}$, where, $M_{\rm Pl}$ is the reduced Planck mass. Thus, if we wish our perturbative calculations to be valid and hence trustworthy during inflation, we would want all the couplings to stay small till super-Planckian scales. This could also help in keeping the inflaton potential sufficiently flat. Thus, we impose the constraint that the maximum value a growing coupling should take must be $\sqrt{4\pi}$ at $\mu = 1.2 \times 10^{20}$ GeV $\approx 50 M_{\rm Pl}$ so that perturbativity (in matter sector) breaks down at scales $\mu \geq 50 M_{\rm Pl}$ only. Of course, at $\mu \sim M_{\rm Pl}$, gravity is expected to become strongly coupled but this is a separate subject in itself.

Let us now impose this constraint: we would like to have

$$g_1(\mu) < \sqrt{4\pi} \quad \text{for } \mu < 1.2 \times 10^{20} \text{ GeV}. \tag{17}$$

Assuming the mass of the fermion to be 1.2 TeV, using the fact that g_1 at the scale 1.2 TeV is 0.47 and using Eq. (14), we find that the corresponding allowed maximum value of Y^2 is

$$Y^2_{\rm max} = \begin{cases} 2.08, & \text{model X}, \\ 1.04, & \text{model Y}. \end{cases} \tag{18}$$

This tells us that if we impose this constraint (and assume $m_f = 1.2$ TeV), then, in model Y, the case with $Y = 7/6$ (i.e. $Q_1 = 5/3$ and $Q_2 = 2/3$) is not consistent with our requirements.[1] Therefore the case which will interest us most will be the one in which the hypercharge of the fermion doublet is chosen to be $-5/6$.

Continuing with the assumption that the mass of the fermion is 1.2 TeV, Eq. (18) must be compared with what one obtains by combining Eq. (2) and Eq. (4)

$$Q^2 = \frac{1}{1.32 y_m} \left(\frac{\sigma}{\rm fb}\right)^{1/2}, \tag{19}$$

where, $Q^2 = Q_1^2 + Q_2^2$ so that

$$Q^2 = 2Y^2 + \frac{1}{2}. \tag{20}$$

In Fig. 2, we plot Q^2 against the allowed range of cross section $\sigma(pp \to \gamma\gamma)$, for various choices of the Yukawa coupling y_m. Imposing the additional requirement of Eq. (18) ensures that for a given value of cross section, there is a unique value of the Yukawa coupling (for m_f fixed at 1.2 TeV).

The most important lesson to be learnt from Fig. 2 is that for Model X, the Yukawa coupling y_M consistent with (i) $\sigma > 3$ fb, (ii) not having Landau pole in g_1 below 10^{20} GeV (assuming $m_f = 1.2$ TeV) is such that $y_M > 0.63$. Similarly, for model Y, the Yukawa coupling consistent with these requirements is $y_M > 0.5$. Since the Yukawa coupling cannot be kept arbitrarily small, maintaining perturbativity till Planck scale is a fairly nontrivial constraint.

In summary, if Eq. (4) was the only constraint, then we could have decreased y_M as much as we liked by increasing the quantity Q^2 but this quantity is fixed by Eq. (18) and so we no longer have that freedom. Second, if some of the values of the Yukawa coupling in this allowed range are

[1] Of course, one could change the mass of the fermion and could argue that hitting Landau pole around Planck scale is not undesirable. In that case one could work with $Y = 7/6$ fermion as is done in [110].

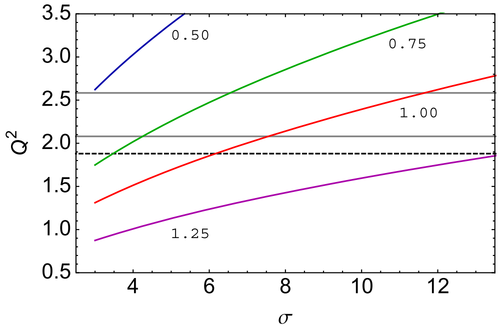

FIG. 2. The uppermost horizontal line here is the value of Q^2 one gets by imposing the requirement that g_1 is less than $\sqrt{4\pi}$ for scales below 10^{20} GeV (see text for details) for the mass of the fermion fixed to be 1.2 TeV in model Y. The middle horizontal line is the value of Q^2 one gets by imposing the same requirement in model X. The lowest horizontal line is the value of Q^2 for the hypercharge $Y = -5/6$ for the multiplet in model Y. The curves correspond to the Q^2 found from Eq. (4) for the values of the Yukawa coupling y_m shown. For a given $\sigma(pp \to S \to \gamma\gamma)$, this means that there is a unique allowed value of the Yukawa coupling y_m.

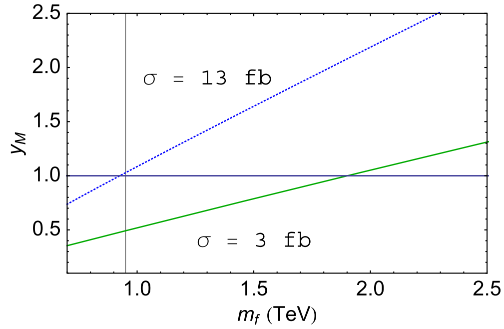

FIG. 3. The inferred value of Yukawa coupling for Model Y at the scale m_f for any chosen m_f (the mass of the fermion) and for the two extreme values of the diphoton rate obtained with hypercharge of the fermion doublet chosen to be $-5/6$ i.e. within the limits specified by Eq. (18). We find that having lighter fermion helps in keeping the coupling small. Though for values of m_f below 950 GeV (i.e. left of the vertical line), all cross sections of the diphoton signal can be explained, these values are below the lower limit on the mass of vectorlike fermion of charge -4/3 obtained by the LHC (see [105]).

somehow not viable, then there is no scope for explaining the corresponding value of cross-section from this model while maintaining perturbativity and $m_f = 1.2$ TeV.

Let us now see what happens when we change the mass of the fermions in the doublet $(\psi_{-1/3}\psi_{-4/3})^T$ considered in model Y. As we saw in Sec. II [in the discussion just before Eq. (4)], the experimental lower limit on the mass of vectorlike fermions of charge $-4/3$ has been obtained by the ATLAS collaboration and is around 0.95 TeV [105]. Using Eq. (2) and Eq. (3) along with the requirement imposed by Eq. (17) ensures that in this scenario, increasing the mass of fermion increases the inferred value of Yukawa coupling of the fermion in order to match the observed diphoton rate. This is illustrated in Fig. 3 which also shows that if $m_f > 1.7$ TeV, no allowed value of diphoton rate can be explained while keeping the Yukawa coupling smaller than one. From the RG equation of λ_S (see appendix), it is worth noticing that having a value of $y_M > 1$ can be an impediment in maintaining perturbativity till Planck scale since in such a case λ_S and λ_{SH} can undergo huge RG evolution. From Fig. 3, we also learn that if $m_f < 0.9$ TeV, all the allowed values of the diphoton rate can be explained with Yukawa couplings smaller than unity. However, this small value of m_f is not favored by the ATLAS results [105].

IV. PERTURBATIVITY, VACUUM STABILITY AND OTHER CONSTRAINTS

In order to analyse the issues of perturbativity and vacuum stability, we need to find the renormalization group (RG) evolution equations for the scenario we are considering. The one-loop RG evolution equations for a model with SM particles, a real singlet scalar and $SU(2)_L$ singlet vectorlike quark (i.e. $Q = 2/3$) have already been found in Ref. [107]. In order to deal with model X and model Y, we need to generalize the equations presented in [107] to take into account both $SU(2)_L$ singlet and $SU(2)$ doublet vectorlike fermions with exotic charges i.e. $Q > 2/3$.

A. Solution of RGEs: Perturbativity

The detailed RG evolution equations for both model X and model Y have been given in Appendix A. We numerically solve the RG equations for three different regimes: (i) from $m_t(= 173.1$ GeV) to $m_S(= 750$ GeV), (ii) from m_S to m_f (the mass of the fermion), and, (iii) from m_f to Planck scale and beyond. In regime (i), one works with the RG equations of the SM (g_1, g_2, g_3, y_t and $\lambda_H^{\rm SM}$) while in regime (ii), two new scalar couplings (λ_S and λ_{SH}) enter the description. The threshold corrections due to the existence of the singlet scalar S with mass 750 GeV causes the familiar discontinuity in the evolution of λ_H [109] as given in Eq. (9). Finally, in regime (iii), in model X, we have Yukawa coupling y_M and in model Y, we have Yukawa couplings y_M and y_F.

In Fig. 4, we have shown the RG evolution of the scalar couplings λ_H, λ_S and λ_{SH}, the gauge coupling g_1 and the Yukawa coupling y_M for model X for a choice of parameters mentioned in the figure caption. This particular choice of parameters can easily explain the cross section $\sigma \approx 4$ fb using model X. Similarly, in Fig. 5, we have shown the RG evolution of the scalar couplings λ_H, λ_S and λ_{SH}, the gauge coupling g_1 and the Yukawa couplings y_M and y_F for model Y for the choice of parameters mentioned in the figure caption. This particular choice of parameters

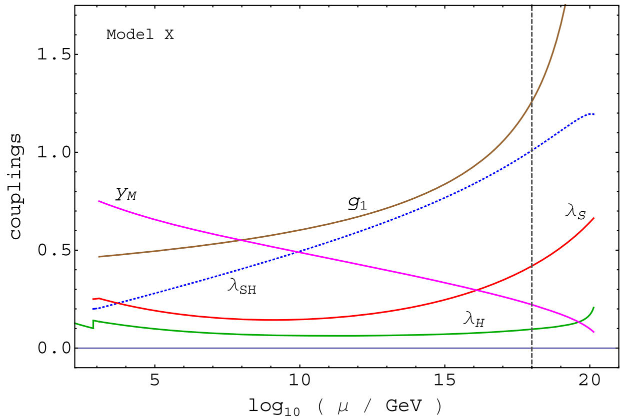

FIG. 4. For model X: the evolution of the various couplings for the case $y_M(m_f) = 0.75$, $\lambda_S(m_s) = 0.25$ and $\lambda_{SH} = 0.20$. Here $m_f = 1.2$ TeV, and this choice of parameters can explain the diphoton excess if $\sigma(pp \to S \to \gamma\gamma) \approx 4$ fb. The square of the electric charge (and hence hypercharge) of the fermion is chosen to be $Q^2 = Q^2_{\rm max} = 2.08$ [see Eq. (18)]. Notice that g_1 is below $\sqrt{4\pi}$ for $\mu < 1.2 \times 10^{20}$ GeV. The region to the right of the vertical line corresponds to $\mu > 10^{18}$ GeV where quantum gravitational effects are expected to be important.

can easily explain the cross section $\sigma \approx 3$ fb using model Y. It is noteworthy that, as can be guessed from Fig. 3, all observed values of the diphoton rate can be explained by considering a light enough fermion while maintaining perturbativity till $\mu \sim 50 M_{\rm Pl}$ by appropriately choosing the λ_S and λ_{SH} at the scale of mass of the singlet scalar. Also, as is clear from Fig. 4 and Fig. 5, we evolve the RG equations till well beyond Planck scale. It is noteworthy that the couplings λ_H, λ_S and λ_{SH} are such that not only is

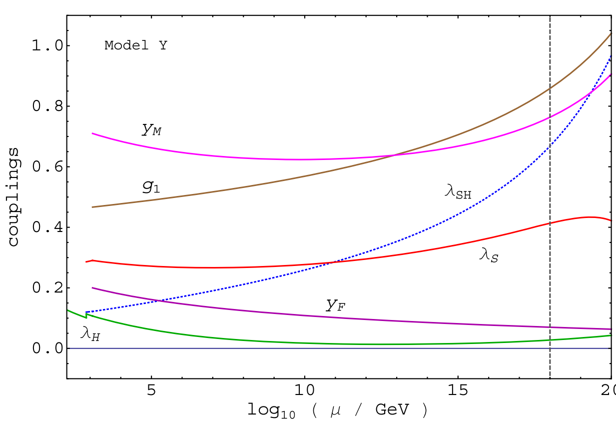

FIG. 5. For model Y: the evolution of the various couplings for the case $y_M(m_f) = 0.71$, $\lambda_S(m_s) = 0.285$ and $\lambda_{SH} = 0.12$. Here $m_f = 1.2$ TeV, and this choice of parameters can explain the diphoton excess if $\sigma(pp \to S \to \gamma\gamma) \approx 3$ fb. The hypercharge of the fermion doublet is chosen to be $Q = -5/6$ [see Eq. (18)]. Notice that g_1 is below $\sqrt{4\pi}$ for $\mu < 1.2 \times 10^{20}$ GeV. The region to the right of the vertical line corresponds to $\mu > 10^{18}$ GeV where quantum gravitational effects are expected to be important.

the mixing of S with SM Higgs H negligible, but also, the mass eigenvalues take up the correct values m_S and m_H.

On the other hand, what we have shown is that both model X and model Y stay weakly coupled (since all the couplings stay smaller than unity) all the way till Planck scale while still explaining the observed diphoton excess.

Perturbativity issue with multiple generations of up-type vector-like fermion: Let us discuss the situation if instead of considering vectorlike fermions with exotic charges, we consider vectorlike fermions with up-type quark charge but increase the number of flavors/generations of the fermions. If we consider N_f vectorlike fermions, the Eq. (14) will be modified to

$$b_1 = \begin{cases} \frac{41}{10} + \frac{12}{5} N_f Y_f^2, & \text{model X,} \\ \frac{41}{10} + \frac{24}{5} N_f Y_f^2, & \text{model Y,} \end{cases} \tag{21}$$

where Y_f represents the hypercharge of vectorlike fermion. Similarly, the equation (18) representing the maximum allowed value of hypercharge from the Landau pole requirement will be modified to.

$$(N_f Y_f^2)^{\rm max} \sim \begin{cases} 2.08 & \text{model X,} \\ 1.04 & \text{model Y.} \end{cases} \tag{22}$$

Incorporating $Y_f = 2/3$ for a $SU(2)$ singlet up-type vector-like fermion in model X and $Y_f = 1/6$ for a $SU(2)$ doublet up-type vectorlike fermion in model Y, we obtain

$$N_f^{\rm max} \sim \begin{cases} 4 & \text{model X,} \\ 37 & \text{model Y.} \end{cases} \tag{23}$$

We analyze two cases of model X in which one can explain at least the minimum value of diphoton cross section [$\sigma(pp \to S \to \gamma\gamma) \approx 3$ fb] by considering multiple up-type vectorlike fermions and particular value of Yukawa couplings. In those cases, we consider (i) $y_M = 1.5$, $N_f = 2$ and (ii) $y_M = 0.75$, $N_f = 4$. As is clear from Fig. 6, the Yukawa coupling becomes large at a very low energy scale in both the cases. The reason behind this is the competition between the positive contribution coming from the term of the type $N_f y_M^2$ and the negative contribution coming from the term of the type $Q_f g_1^2$ in the RG evolution equation of the Yukawa coupling y_M [see Eq. (A14) in the Appendix). If the former gets dominant (due to large N_f), it will make the theory nonperturbative at a low energy scale. However, if the latter become dominant by increasing value of Q_f, it keeps the theory perturbative till Planck scale. Therefore, we show that though increasing number of generation does help in explaining the diphoton cross-section at TeV scale (see Sec. II), it does not allow us to keep the theory perturbative all the way till Planck scale.

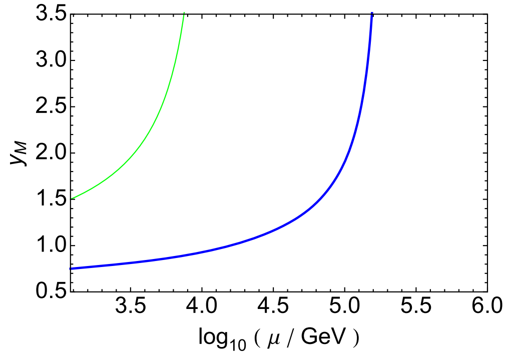

FIG. 6. RG evolution of Yukawa coupling y_M for two cases with multiple copies (N) of vectorlike fermion in model X: (i) The solid (green) line shows the evolution of y_M for the parameters $y_M(m_f) = 1.5$, $N = 2$, $Q = 2/3$ and $\sigma(pp \to S \to \gamma\gamma) \approx$ 3 fb, (ii) The thick (blue) line shows the evolution of y_M for the parameters $y_M(m_f) = 0.75$, $N = 4$, $Q = 2/3$ and $\sigma(pp \to S \to \gamma\gamma) \approx 3$ fb.

B. Vacuum stability and inflation

In the last section, we found two models compatible with the observed diphoton excess and a regime in which these models stay weakly coupled till Planck scale. Let us now find a few more consequences of this. It is well known that at energy scale above 10^{10} GeV, the self-coupling of the SM Higgs turns negative [111–114]. This observation has motivated many scenarios of physics beyond the standard model. It is well known that adding more fermions typically causes λ_H to turn negative at even smaller scales. On the other hand, addition of more scalars can potentially enhance λ_H so much that it may hit Landau pole below Planck scale.

Before we discuss electroweak vacuum stability in model X and model Y, it is worthwhile recalling the conditions of stability. At large field values, the scalar potential is of the form

$$V = \lambda_H(\Phi^\dagger\Phi)^2 + \lambda_S\chi^4 + \lambda_{SH}\chi^2(\Phi^\dagger\Phi). \tag{24}$$

From this, it is clear that if all the three couplings λ_H, λ_S and λ_{HS} are positive, the potential will be positive and hence stable (since the SM vacuum corresponds to $V = 0$). Had the term λ_{HS} been negative, the requirement of obtaining a positive scalar potential for asymptotically large values of the fields leads to the conditions: $4\lambda_S\lambda_H > \lambda_{SH}^2$, $\lambda_H > 0$, $\lambda_S > 0$.

As can be seen in Fig. 4 and Fig. 5, in the present scenario, the self coupling of the SM Higgs (λ_H) stays positive and small till Planck scale. Moreover, the self coupling of the real singlet scalar λ_S and the coupling between the Higgs and the real singlet scalar (λ_{SH}) also stay positive and small till Planck scale. Thus, in these models, with the choice of parameters we have shown, the vacuum stability problem is solved. At this stage, we must mention that the instability scale in SM is 10^8 GeV at one-loop accuracy, 10^{10} GeV at two-loop accuracy and even higher at three-loop accuracy [113]. In our calculations, we have restricted our analysis to one-loop accuracy. Since the instability scale gets pushed up in energy as one increases the accuracy of calculation, and since the vacuum gets stabilized by the threshold effects [109] of the newly added scalar (at 750 GeV), one expects the stability to be maintained if one repeats these calculations at two-loops. Still, it may be worthwhile to check our conclusions for two-loop accuracy.

The recent combined results from Planck, BICEP and KECK [115,116] experiments tell us that inflation with both $\lambda\phi^4$ potential and $m^2\phi^2$ potential are ruled out by upper limits on primordial B-mode polarization of CMB. Since this is typically the form of the potential at large field values for most particle physics scenarios, this is a very important observational constraint. A well known way out is to consider nonminimal coupling of the inflaton to gravity of the form $\xi R\phi^2$ where R is the Ricci scalar. This nonminimal coupling is actually inevitable from the point of view of supergravity. It is well known that such scenarios can easily be made compatible with data.

In SM, the negativity of the self coupling λ_H at high scale makes it difficult to use SM Higgs as the inflaton even with a nonminimal coupling to gravity (see however [117] for a recent attempt). In our model, at high scales, the potential in both H direction and S direction is of quartic form with $\lambda \sim \mathcal{O}(0.1) > 0$. It is then obvious that a nonminimal coupling of these fields to gravity of the form $\xi R\phi^2$ with $\xi \sim 10^4$ (from Eq. (19) of [118]) leads to successful inflation. Assuming the number of e-foldings of inflation to be 57.7, one obtains (to lowest order in $1/\xi$),

$$n_s \approx 0.96 \quad \text{and} \quad r \approx 0.003, \tag{25}$$

which are compatible with the observations of Planck. Notice that since the corresponding nonminimal coupling is large, one may need to worry about quantum corrections such as those dealt with in [119]. Given this, one should reanalyze the RG evolution of all couplings after including the effects of the nonminimal coupling ξ (see e.g. [120,121], where this issue is partially addressed). We leave this kind of analysis for future work.

Thus, we have shown that (i) the vacuum stability problem gets solved in both model X and model Y, and, (ii) the both the SM Higgs and the 750 GeV real singlet scalar whose decay causes the diphoton excess are very good candidates for inflaton in these models.

C. Additional constraints

Let us now consider an additional constraint which distinguishes between the two models we have been working with. In model X, in order to explain the diphoton excess and maintain perturbativity, we chose the charge of the fermion to be exotic. But there is no decay mode for a particle with this value of charge. This is because, as can be seen from the discussion in Sec. III B, the only way a model X fermion can decay into SM particles is if its charge is either $2/3$ or $-1/3$. If we consider the charge to be arbitrary, it will be stable and may have a dangerous relic abundance.

Given that the fermions are $SU(3)$ charged as well, one may think that the fermions annihilate to negligible abundances because of the strong interactions. For the mass of fermion $m_f \sim$ TeV, the relic abundance turns out to around 10^{-4} per nucleon. Though this is quite small, according to [122], the abundance of any particles with charge more than 0.83 must be less than 5×10^{-22} per nucleon. Therefore, it may seem that model X is not viable.

However, since the fermions are colored, the story could be more complicated. As mentioned in [123], for $T < \Lambda_{\mathrm{QCD}}$, the colored heavy fermions can form bound states ("pions") made up of these exotic fermions and effectively have an enhanced cross section for annihilation (given by the capture cross section) so that all fermions which form bound states annihilate away. Thus, the relic abundance is given by the fraction of these fermions which did not become bound. Therefore, the overall relic abundance of exotic fermions might be less than the one found above (i.e. 10^{-4}). In this sense the concern about the longevity of the fermion in model X is similar to that of the long-lived gluinos in split-SUSY [124], we do not address this issue here.

Instead we argue that the $SU(2)_L$ doublet fermion in model Y decays into SM particles. In model Y, since the fermion is $SU(2)_L$ doublet, the Lagrangian contains terms such as

$$\mathcal{L} = \mathcal{L}^{\mathrm{SM}}_{\mathrm{Yuk}} + \mathcal{L}^{\mathrm{SM}}_{\mathrm{gauge}} + y_M \chi \bar{\psi}_L \psi_R + y_F \Phi^c \bar{\psi}_L d_R + g_2 \bar{\psi}_2 \gamma^\mu W^-_\mu \psi_1 + g_2 \bar{\psi}_1 \gamma^\mu W^+_\mu \psi_2. \quad (26)$$

Because of the presence of these terms, ψ_1 decays into a bottom quark and Higgs while ψ_2 decays to W^-, bottom quark and Higgs. Hence from this point of view, model Y is the preferred model which serves our purpose.

Notice that this requirement of decay of the fermion can be used to rule out not only model X but also all values of the hypercharge in model Y other than $7/6$ and $-5/6$. But since the hypercharge $7/6$ leads to a Landau pole in g_1 around Planck scale, we prefer to work with model Y with hypercharge of the fermion doublet chosen to be -5/6.

V. SUMMARY

One of the most peculiar features of the SM is that it stays perturbative and hence calculable till Planck scale. This does not hold well for arbitrary variations of the SM. In this work, we first began by arguing that if the 750 GeV diphoton excess is interpreted as being due to heavy scalar (heavier cousin of the SM Higgs) coupled to the photons through a fermion loop, then it is generically difficult to maintain perturbativity as we probe higher energies. We did this by demonstrating that the inferred values of the couplings in any scenario which explains the diphoton excess are generically large (see Fig. 1).

We then presented a scenario which requires adding very few new particles and which stays weakly-coupled till Planck scale. These new particles are a singlet real scalar (whose decay causes diphoton excess) with mass 750 GeV and a colored vectorlike fermion with exotic charge $Q > 2/3$ and mass m_f such that $m_f \sim 1.2$ TeV. We studied two variations of this scenario: in model X, the newly added fermion is a singlet under $SU(2)_L$ while in model Y, the newly added fermion is a doublet under $SU(2)_L$. We argued that the charge of the newly added colored fermion can not be that of a quark (i.e. $2/3$ or $1/3$) in order to explain the observed diphoton excess if one wishes to maintain perturbativity till high scale. We then limited the possible hypercharges of the fermions in the two models from the requirement that the gauge coupling g_1 stays sufficiently small till a little above Planck scale (see Fig. 2). Moreover, we chose the couplings of the scalar sector such that the newly added scalar has negligible mixing with the SM Higgs and the mass eigenvalues take up the correct values of the masses of the scalar particles.

We then showed that using a lighter fermion ($m_f \sim 900$ GeV), one can (i) explain the diphoton excess over the entire range of σ[3–13] fb, and also, (ii) maintain perturbativity till Planck scale. Using a heavier fermion ($m_f \sim 1200$ GeV), it is still possible to maintain perturbativity till Planck scale for some choice of parameters but it becomes difficult to explain all the possible values of the diphoton production cross section (in this context see Fig. 3). However, as we discussed in Sec. II, the experimental lower limits on the masses of any new vectorlike fermions (with various charges) are around 0.85–0.95 TeV [103–105], so, we mostly considered the mass of the fermions to be bigger than that, around 1.2 TeV. From Fig. 3, this value of fermion mass successfully explains only smaller value of the cross-section i.e. [3–9] fb.

We found the renormalization group evolution equations for these models and evolved the couplings with this particle content and reasonable choices of parameters. This helped us in finding a regime in which these models stay weakly coupled till Planck scale, the scalar couplings λ_H, λ_S and λ_{SH} stay positive i.e. the vacuum stability problem is solved and one obtains successful inflation. We showed that if instead of adding fermions with a large value of charge, we add multiple copies of a fermion with charge $2/3$, we are unable to maintain perturbativity all the way till Planck scale.

We then argued that in model X, since the charge of the fermion is large, it does not decay into any SM particles and it can lead to dangerous abundance of the newly added fermions in the Universe. Thus, the proposed model Y is the only one satisfying all the constraints we have imposed and is hence the preferred model. Since the electric charges of the fermion doublet in model Y are -4/3 and -1/3, the recent observational limits of [105] help us put a lower limit on the mass of the fermions forming the doublet.

ACKNOWLEDGMENTS

The authors would like to thank S. Mohanty and N. Mahajan (PRL, Ahmedabad) for discussions. The authors would also like to thank the anonymous referee for useful comments and suggestions.

APPENDIX: RENORMALIZATION GROUP EVOLUTION EQUATIONS

In this appendix, we present the renormalization group evolution equations (RGEs) for our scenario.

1. Modified RGEs for gauge couplings:

The RG evolution equation for a coupling is found from its beta function, for a gauge coupling g,

$$\frac{dg}{dt} = \beta(g), \tag{A1}$$

where $t = \log(\mu/M_t)$, μ being the energy scale of interest and M_t being the mass of top quark.

The one-loop beta function for a general $SU(N)$ gauge theory is given by (see Eq. (73.41) and Eq. (78.36) of [125] and Eq. (6.38) of [126])

$$\beta(g) = -\frac{1}{3}\frac{g^3}{(4\pi)^2}(11T(A) - 4T(R_{DF}) - T(R_{CS})) + \mathcal{O}(g^2), \tag{A2}$$

where $T(A)$ is the index of the adjoint representation of gauge group, $T(R_{DF})$ is the index of the representation of any Dirac fermion in the theory while $T(R_{CS})$ is the index of the representation of any complex scalar in the theory. Recall that if we have Weyl (or Majorana) fermions in theory, they contribute half as much as a Dirac fermion. Similarly, if we have real scalars in theory, they contribute half as much as a complex scalar. Given this, it is clear that the RG equation for g_3 in model X is

$$\frac{dg_3}{dt} = \frac{g_3^3}{(4\pi)^2}\left(-7 + \frac{2}{3}N_f\right), \tag{A3}$$

while the RG equation for g_3 in model Y is

$$\frac{dg_3}{dt} = \frac{g_3^3}{(4\pi)^2}\left(-7 + \frac{4}{3}N_f\right), \tag{A4}$$

where -7 in the braces is the contribution from SM particles alone. There is an extra contribution from N_f vectorlike fermions but none from the scalar since it is a gauge singlet. In model X, since none of the extra particles are charged under $SU(2)_L$ of the SM, the RG equation for g_2 remains unchanged i.e.

$$\frac{dg_2}{dt} = \frac{g_2^3}{(4\pi)^2}\left(-\frac{19}{6}\right). \tag{A5}$$

For model Y, the RG equation for g_2 is

$$\frac{dg_2}{dt} = \frac{g_2^3}{(4\pi)^2}\left(-\frac{19}{6} + \frac{12}{6}N_f\right). \tag{A6}$$

Similarly, the one-loop beta function for a $U(1)$ gauge theory is given by (see Eq. (66.29) of [125])

$$\beta(g) = \frac{g^3}{(4\pi)^2}\left(\sum_\psi Q_\psi^2 + \frac{1}{4}\sum_\phi Q_\phi^2\right), \tag{A7}$$

where Q_ψ is the $U(1)$ charge of a Dirac fermion and Q_ϕ is the $U(1)$ charge of a complex scalar. Using this, we can find the RG equation for g_1 in model X[2]

$$\frac{dg_1}{dt} = \frac{g_1^3}{(4\pi)^2}\left(\frac{41}{10} + \frac{12N_f Q_f^2}{5}\right). \tag{A8}$$

The RG equation for g_1 in model Y is

$$\frac{dg_1}{dt} = \frac{g_1^3}{(4\pi)^2}\left(\frac{41}{10} + \frac{24N_f Q_f^2}{5}\right). \tag{A9}$$

As previously, the first term in the parenthesis viz $41/10$ is the contribution of SM field content and the second term is due to the extra N_f Dirac fermions of hypercharge Q_f.

2. Modified RGEs for Yukawa couplings:

For a Yukawa coupling, the beta function is defined by

$$\frac{dy}{dt} = \frac{y}{(4\pi)^2}\beta(y). \tag{A10}$$

In SM, the Yukawa couplings take the form of three complex 3×3 matrices, Y_u, Y_d, Y_e (where, the matrix Y_u is responsible for giving mass to the up type quarks of all the generations etc). The one-loop beta functions of Yukawa coupling matrices for SM are given by (see Eq. (6.42) of [126])

[2]Notice that, as is always done, we replace g_1 by $\sqrt{\frac{3}{5}}g_1$ in order to be consistent with grand unified theory (GUT) normalization.

$$\beta_u^{\rm SM} = \frac{3}{2}(Y_u^\dagger Y_u - Y_d^\dagger Y_d) + T - \left(\frac{17}{20}g_1^2 + \frac{9}{4}g_2^2 + 8g_3^2\right),$$
$$\beta_d^{\rm SM} = \frac{3}{2}(Y_d^\dagger Y_d - Y_u^\dagger Y_u) + T - \left(\frac{1}{4}g_1^2 + \frac{9}{4}g_2^2 + 8g_3^2\right),$$
$$\beta_e^{\rm SM} = \frac{3}{2}(Y_e^\dagger Y_e) + T - \frac{9}{4}(g_1^2 + g_2^2), \tag{A11}$$

where

$$T = \mathrm{Tr}(3Y_u^\dagger Y_u + 3Y_d^\dagger Y_d + Y_e^\dagger Y_e). \tag{A12}$$

The one-loop correction to the Yukawa coupling comes from one-loop graphs which contribute to scalar-fermion-fermion vertex. As is beautifully explained in [126], among these, the contribution T comes from the diagrams involving a fermion loop to the scalar propagator while the first set of terms in the above equations come from a Higgs loop correction to the vertex. The last set of terms come from the loops involving gauge fields.

In the scenario we are dealing with, we add N_f vectorlike quarks having charge Q_f and the SM singlet S. In model X, we shall be interested in two cases: (a) $Q_f = 2/3$ and (b) $Q_f > 2/3$. Recall that when $Q_f = 2/3$, the vectorlike fermion couples to SM quarks and Higgs with the Yukawa coupling y_T while when $Q_f > 2/3$, the Yukawa coupling y_T is zero. It is easy to generalize the RG equations presented in [107], we find that

$$\beta_u = \beta_u^{\rm SM} + \sum_{i=1}^{N_f}\left(\frac{3}{2}Y_T^{(i)\dagger}Y_T^{(i)} + 3Y_T^{(i)\dagger}Y_T^{(i)}\right),$$
$$\beta_d = \beta_d^{\rm SM} - \sum_{i=1}^{N_f}\left(\frac{3}{2}Y_T^{(i)\dagger}Y_T^{(i)} + 3Y_T^{(i)\dagger}Y_T^{(i)}\right),$$
$$\beta_e = \beta_e^{\rm SM}. \tag{A13}$$

Moreover, the beta functions of the extra Yukawa couplings in Model X are of the form (see also [94])

$$\beta_{y_T}^{(k)} = \sum_{i=1}^{N_f}\frac{9}{2}Y_T^{(i)\dagger}Y_T^{(i)} + \frac{Y_M^{(k)2}}{2} + \frac{3}{2}(Y_u^\dagger Y_u - Y_d^\dagger Y_d) + T - \left(\frac{17}{20}g_1^2 + \frac{9}{4}g_2^2 + 8g_3^2\right),$$
$$\beta_{y_M}^{(k)} = \sum_{i=1}^{N_f}2\left(\frac{3}{2}Y_M^{(i)2} + 3Y_M^{(i)2}\right) + Y_T^{(k)\dagger}Y_T^{(k)} - \left(\frac{9}{10}(Y_f)^2 g_1^2 + 8g_3^2\right), \tag{A14}$$

where Y_f is the hypercharge of vectorlike fermions. For comparison, one could refer to the appendix of [107] to find the corresponding result in the case of a single vectorlike fermion. Similarly, the beta functions of the extra Yukawa couplings in Model Y can be worked out and are of the form

$$\beta_{y_F}^{(k)} = \sum_{i=1}^{N_f}\frac{9}{2}Y_F^{(i)\dagger}Y_F^{(i)} + \frac{Y_M^{(k)2}}{2} + 3(Y_d^\dagger Y_d) + \mathrm{Tr}(3Y_u^\dagger Y_u + 3Y_d^\dagger Y_d + Y_e^\dagger Y_e) - \left(\frac{17}{20}g_1^2 + \frac{9}{4}g_2^2 + 8g_3^2\right),$$
$$\beta_{y_M}^{(k)} = \sum_{i=1}^{N_f}2\left(\frac{3}{2}Y_M^{(i)2} + 4Y_M^{(i)2}\right) + \frac{1}{2}Y_F^{(k)\dagger}Y_F^{(k)} - \left(\frac{36}{10}(Y_f)^2 g_1^2 + \frac{9}{2}g_2^2 + 8g_3^2\right), \tag{A15}$$

3. Modified RGEs for scalar couplings

The scalar beta functions are related to RG equations by

$$\frac{d\lambda}{dt} = \frac{1}{(4\pi)^2}\beta_\lambda. \tag{A16}$$

In SM, the beta function of the Higgs self-coupling is given by (see Eq. (6.48) of [126])

$$\beta_\lambda^{\rm SM} = 12\lambda^2 - \left(\frac{9}{5}g_1^2 + 9g_2^2\right)\lambda + \frac{9}{4}\left(\frac{3}{25}g_1^4 + \frac{2}{5}g_1^2g_2^2 + g_2^4\right) + 4T\lambda - 4H, \tag{A17}$$

where T is as defined in Eq. (A12) and H is given by

$$H = \mathrm{Tr}(3(Y_u^\dagger Y_u)^2 + 3(Y_d^\dagger Y_d)^2 + (Y_e^\dagger Y_e)^2). \tag{A18}$$

In model X, the couplings λ_H, λ_S and λ_{SH} evolve in accordance with (see [107])

$$\frac{d\lambda_H}{dt} = \frac{2}{(4\pi)^2}\Bigg(12\lambda_H^2 + 6y_t^2\lambda_H - \frac{9}{10}g_1^2\lambda_H - \frac{9}{2}g_2^2\lambda_H + \frac{27}{400}g_1^4 + \frac{9}{16}g_2^4 + \frac{9}{40}g_1^2g_2^2 - 3y_t^4 + \frac{\lambda_{SH}^2}{4} + \sum_{i=1}^{N_f}6\lambda_H y_T^{(i)2} + \sum_{i,j=1}^{N_f}(-3)y_T^{(i)2}y_T^{(j)2} + \sum_{i=1}^{N_f}(-6)y_T^{(i)2}y_t^2\Bigg),$$

$$\frac{d\lambda_{SH}}{dt} = \frac{2}{(4\pi)^2}\Bigg(\lambda_{SH}\Bigg(2\lambda_{SH} + 6\lambda_H + 3y_t^2 - \frac{9}{20}g_1^2 - \frac{9}{4}g_2^2 + 3\lambda_S + \sum_{i=1}^{N_f}6y_M^{(i)2} + \sum_{i=1}^{N_f}3y_T^{(i)2}\Bigg) + \sum_{i=1}^{N_f}(-12)y_T^{(i)2}y_M^{(i)2}\Bigg),$$

$$\frac{d\lambda_S}{dt} = \frac{2}{(4\pi)^2}\Bigg(9\lambda_S^2 + \lambda_{SH}^2 + \sum_{i=1}^{N_f}12\lambda_S y_M^{(i)2} + \sum_{i=1}^{N_f}(-12)y_M^{(i)4}\Bigg). \tag{A19}$$

In model Y, the couplings λ_H, λ_S and λ_{SH} evolve in accordance with

$$\frac{d\lambda_H}{dt} = \frac{2}{(4\pi)^2}\Bigg(12\lambda_H^2 + 6y_t^2\lambda_H - \frac{9}{10}g_1^2\lambda_H - \frac{9}{2}g_2^2\lambda_H + \frac{27}{400}g_1^4 + \frac{9}{16}g_2^4 + \frac{9}{40}g_1^2g_2^2 - 3y_t^4 + \frac{\lambda_{SH}^2}{4} + \sum_{i=1}^{N_f}6\lambda_H y_F^{(i)2} + \sum_{i,j=1}^{N_f}(-3)y_F^{(i)2}y_F^{(j)2} + \sum_{i=1}^{N_f}(-6)y_F^{(i)2}y_b^2\Bigg),$$

$$\frac{d\lambda_{SH}}{dt} = \frac{2}{(4\pi)^2}\Bigg(\lambda_{SH}\Bigg(2\lambda_{SH} + 6\lambda_H + 3y_t^2 - \frac{9}{20}g_1^2 - \frac{9}{4}g_2^2 + 3\lambda_S + \sum_{i=1}^{N_f}12y_M^{(i)2} + \sum_{i=1}^{N_f}3y_F^{(i)2}\Bigg) + \sum_{i=1}^{N_f}(-24)y_F^{(i)2}y_M^{(i)2}\Bigg),$$

$$\frac{d\lambda_S}{dt} = \frac{2}{(4\pi)^2}\Bigg(9\lambda_S^2 + \lambda_{SH}^2 + \sum_{i=1}^{N_f}24\lambda_S y_M^{(i)2} + \sum_{i=1}^{N_f}(-24)y_M^{(i)4}\Bigg). \tag{A20}$$

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