Diphoton excess at 750 GeV in leptophobic U(1)′ model inspired by E_6 GUT

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ABSTRACT: We discuss the 750 GeV diphoton excess at the LHC@13TeV in the framework of leptophobic U(1)′ model inspired by E_6 grand unified theory (GUT). In this model, the Standard Model (SM) chiral fermions carry charges under extra U(1)′ gauge symmetry which is spontaneously broken by a U(1)′-charged singlet scalar (Φ). In addition, extra quarks and leptons are introduced to achieve the anomaly-free conditions, which is a natural consequence of the assumed E_6 GUT. These new fermions are vectorlike under the SM gauge group but chiral under new U(1)′, and their masses come entirely from the nonzero vacuum expectation value of Φ through the Yukawa interactions. Then, the CP-even scalar h_Φ from Φ can be produced at the LHC by the gluon fusion and decay to the diphoton via the one-loop diagram involving the extra quarks and leptons, and can be identified as the origin of diphoton excess at 750 GeV. In this model, h_Φ can decay into a pair of dark matter particles as well as a pair of scalar bosons, thereby a few tens of the decay width may be possible.

KEYWORDS: Beyond Standard Model, Higgs Physics

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Contents

1 Introduction

The LHC Run-II experiment started taking data at the 13 TeV center mass of energy in 2015, and searching new physics signals. One of the promising candidates for new physics is an extra heavy scalar boson, which is predicted by many Beyond Standard Models (BSMs). Therefore the search for such a heavy resonance plays a crucial role in testing BSMs.

Recently, both ATLAS and CMS collaborations reported excesses in the diphoton resonance search around 750 GeV [1, 2]. The local (global) significances are $3.6\,(2.0)\,\sigma$ (ATLAS) and $2.6\,(1.2)\,\sigma$ (CMS), respectively. An interesting point is that the best fit value of the decay width of the resonance in the ATLAS data is ~ 45 GeV, which is about 6% of the resonance mass, while the CMS data prefer a narrow decay width [1, 2]. This issue is not still conclusive, but it may be important for theorists to survey the possibility of the diphoton resonance in BSMs [3–130].

One of the attractive candidates for the BSMs which predict the heavy scalar resonance is U(1)′ extension of the Standard Model (SM). In a minimal setup, an additional gauged U(1)′ is introduced and spontaneously broken by the nonzero vacuum expectation value (VEV) of an extra U(1)′-charged scalar (Φ). In the effective lagrangian, an extra massive gauge boson is predicted and a CP-even scalar also appears around the U(1)′ breaking scale. If the scalar mixes with the SM Higgs, it will decay to diphoton, although the scalar mixing is strictly constrained by the Higgs signal strengths and the other resonance searches at the collider experiments. Besides, U(1)′ may cause the gauge anomaly depending on the U(1)′ charge assignments to the SM fermions. Only anomaly-free and generation-independent U(1)′ charge assignment is the linear combination of the hyper-charge and U(1)$_{B-L}$, introducing the right-handed neutrinos. However, such U(1)′ faces a strong bound from the Drell-Yan process, so that the U(1)′ breaking scale can not be low.

As another possibility, we can discuss a U(1)′ model, where the anomaly-free conditions are satisfied by introducing extra chiral fermions. One interesting U(1)′ extension would be the one motivated by the Grand Unified Theory (GUT) with a large rank. For instance, the E_6 GUT principally predicts two additional U(1)′ symmetries, together with extra chiral quarks and leptons. The extra chiral fermions are vector-like under the SM gauge groups but chiral under the two U(1)′. Interestingly, it is suggested that a certain linear combination of the U(1)′ charges could be leptophobic and, in fact, has been widely studied so far in the many literatures [131–142]. Such a leptophobic U(1)′ has been paid an attention to so far, because it can evade the strong bound from the Drell-Yan process. Then the U(1)′ breaking scale can be lower than the U(1)$_{B-L}$ case.

In ref. [140], the present authors discussed the Higgs physics in the leptophobic U(1)′ model in the context of Type-II 2HDM with extra leptophobic U(1)′ instead of softly broken Z_2 symmetry [139]. The reason for 2HDM is that the SM up-type and the down-type fermions carry different U(1)′ charges, so that one can not write Yukawa couplings for all the SM fermions with one Higgs doublet. It is mandatory to extend the Higgs sector by introducing (at least) one more Higgs doublet where two Higgs doublets H_1 and H_2 carry different U(1)′ charges. This way one can write down all the Yukawa couplings for the SM fermions without generating dangerously huge FCNC mediated by neutral Higgs boson [139]. They also found a neutral fermion (ψ_X) in the extra fermions make a good dark matter candidate in this model [140]. One generation of the SM fermions and extra fermions belongs to a single **27** representation of E_6 gauge group, making an anomaly free representation.

In this model, both dark matter and extra chiral fermions get their masses entirely from the nonzero VEV of Φ and the Yukawa couplings. One interesting aspect is that the extra quarks only couple with Φ, but not Higgs doublets because of the U(1)′ charge assignment. Then, the CP-even scalar mode (h_Φ), which appears after the U(1)′ symmetry breaking, does not couple with the SM fermions, neglecting mixing between h_Φ and the other scalars from two Higgs doublets.

Once we assume that the mixing is tiny, we can expect that h_Φ mainly decays through the Yukawa couplings with the extra chiral fermions. As a result, h_Φ can decay to gg and $\gamma\gamma$ via the one-loop diagrams involving the extra chiral fermions which are vector-like under the SM gauge group. In addition, h_Φ can decay into a pair of DM as well as a pair of scalar bosons (hh, Hh, HH, AA) if kinematically allowed, thereby its decay width being increased significantly. This is how we explain the 750 GeV diphoton excess reported by ATLAS and CMS recently. We investigate the parameter region favored by the 750 GeV diphoton excess, and discuss the consistency with the other results on the new physics search at the LHC.

In section 2, we introduce our model setup inspired by E_6 GUT. In section 3, we perform phenomenological analysis on the 750 GeV diphoton excess in our model, and discuss the dark matter physics, according to the interpretation of the diphoton excess. Finally we summarize our results and give some future prospects for probing the scenarios in the future LHC experiments in section 4.

2 Leptophobic U(1)′ model inspired by E_6 GUT

2.1 Model: matter contents and their quantum numbers

Here, we introduce our setup based on refs. [139, 140], where the SM gauge group is extended to include leptophobic U(1)′ gauge symmetry, and additional chiral fermions are introduced in order to cancel gauge anomalies. From the bottom-up point of view, there are many varieties for U(1)′ charge assignments that achieve gauge anomaly cancellations. One simple and interesting way is to consider U(1)′ symmetries predicted by the supersymmetric E_6 GUT, where the all SM fields including Higgs doublets can be unified into three-family **27**-dimensional fields.

Our approach is bottom-up, but let us also explain the underlying theory briefly. The E_6 GUT predicts additional two U(1) symmetries in addition to the SM gauge groups. After the E_6 breaking, one **27**-dimensional fermion field can be decomposed to all the SM fermions in one generation, together with right-handed and left-handed neutrinos, and one extra SM vector-like quark and lepton pairs. If U(1)′, which is given by a linear combination of the two U(1), is assumed to be remained up to the low scale, the anomaly-free condition for the extra U(1)′ can be satisfied by the chiral fermions predicted by the **27**-dimensional fields.

In our model inspired by the E_6 GUT, extra U(1)′ is introduced and its charges are assigned to the SM particles by the linear combination of U(1)$_Y$ and two U(1) of E_6, called U(1)$_\chi$ and U(1)$_\psi$. It is known that one certain linear combination realizes a so-called leptophobic U(1)′, where the SM charged leptons are not charged under the U(1)′ [131–142]. We adopt this leptophobic charge assignments and discuss the phenomenology. The U(1)′ charge assignment for the SM fermions is shown in table 1. The U(1)′ is the subgroup of E_6, so that it is anomaly-free introducing the extra chiral fermions predicted by the **27**-dimensional fields (see for example [143]).

The quantum numbers of the SM and the extra chiral fermions are summarized in table 1. Q^i, d_R^i and u_R^i are the SM quarks, which carry nonzero U(1)′ charges, while L^i and e_R^i are the SM leptons which are not charged under the U(1)′, as shown in table 1. In addition, extra quarks and leptons are contained in the **27**, which are denoted by D_L^i, D_R^i and $\widetilde{H}_L^i$, $\widetilde{H}_R^i$ ($i = 1$, 2, 3). There are left-handed and right-handed neutrinos, n_R^i, N_L^i, whose U(1)′ charges are ± 1. Note that all the matter fields (both fermions and bosons) carry nonzero gauge charges and there are no gauge singlets in this model.

In order to realize electroweak symmetry breaking and U(1)′ breaking, we introduce scalars: two Higgs doublets, denoted by H_1 and H_2, and a U(1)′-charged singlet, Φ. Their quantum numbers are shown in table 1. Note that scalar fields have nothing to do with E_6, since they do not belong to **27**-representation of E_6. We have used the E_6 connection only to the fermion sector in relation with gauge anomaly cancellation. Scalar fields do not contribute to gauge anomaly, and thus need not belong to **27** of E_6 within the bottom-up approach we are taking in this work. As we discuss in the subsection 2.4, H_1 and H_2 are distinguished by their U(1)′ charges and couple with the up-type and the down-type fermions, respectively. Φ develops the nonzero VEV, and breaks U(1)′ gauge symmetry

Fields	spin	SU(3)	SU(2)	U(1)$_Y$	U(1)$'$	Z_2^{ex}
Q^i		**3**	**2**	$1/6$	$-1/3$	
u_R^i		**3**	**1**	$2/3$	$2/3$	
d_R^i		**3**	**1**	$-1/3$	$-1/3$	
L_i	$1/2$	**1**	**2**	$-1/2$	0	$+$
e_R^i		**1**	**1**	-1	0	
n_R^i		**1**	**1**	0	1	
H_2		**1**	**2**	$-1/2$	0	
H_1	0	**1**	**2**	$-1/2$	-1	$+$
Φ		**1**	**1**	0	-1	
D_L^i		**3**	**1**	$-1/3$	$2/3$	
D_R^i		**3**	**1**	$-1/3$	$-1/3$	
$\widetilde{H}_L^i$	$1/2$	**1**	**2**	$-1/2$	0	$-$
$\widetilde{H}_R^i$		**1**	**2**	$-1/2$	-1	
N_L^i		**1**	**1**	0	-1	

Table 1. Matter contents in U(1)$'$ model inspired by E$_6$ GUTs. Here, i denotes the generation index: $i = 1, 2, 3$.

spontaneously. Note that our neutrino is a Dirac particle and the mass term is given by the Yukawa coupling: $y_\nu^{ij}\overline{L_i}H_1 n_{Rj}$.

Note that our model is nonsupersymmetric by construction. Here let us give brief comments on the connection between our nonsupersymmetric U(1)$'$ model and the supersymmetric E_6 GUT, which we might expect to be the underlying theory of our model. In the framework of the supersymmetric E_6 GUT, one **27**-representational chiral superfield realizes all the fermions in one generation shown in table 1. In addition, the bosonic partners with opposite R-parity are predicted. As we can see in table 1, the quantum numbers of H_1 and H_2 are the same as the ones of $\widetilde{H}_L$ and $\widetilde{H}_R$ and the U(1)$'$ charge of Φ is the same as $\overline{n_R^i}$ and N_L^i. In fact, the U(1)$_\chi$ and U(1)$_\psi$ charges of Φ, which are introduced in ref. [140], are the same as the ones of N_L^i. Then, the bosonic fields could be interpreted as the superpartners of the extra fermions, although we have to consider an issue how to decouple the other bosonic fields except $H_{1,2}$ and Φ, which is known to be notoriously difficult in any GUT models.

Another issue may be how to realize the U(1)$'$ in the low energy regime, although our approach is bottom-up. The possibility that the leptophobic U(1)$'$ is generated by kinetic mixing in the supersymmetric E_6 GUT has been studied in refs. [135–137]. Furthermore, Yukawa couplings may cause serious problems in not only E_6 but also SO(10) and SU(5) GUTs. The GUTs unify the matter fields in the elegant ways, but the unification makes it harder to explain the realistic fermion mass matrices.

In this work, our main motivation is to propose one possible setup to explain the diphoton excess, so that we take the bottom-up approach. Our matter contents and U(1)$'$ charge assignment are inspired by the supersymmetric E_6 GUT, but the study about the

explicit consistency with the GUT is clearly beyond our scope. In a sense, our model setup could be considered as a simplified model for E_6 GUT, where fermions form a complete **27**-representation of E_6, but the scalar fields do not. We introduce only a minimal number of scalar fields in order to break the gauge symmetry spontaneously and generate masses for the SM and extra chiral fermions in our model, in such a way that the neutral Higgs-mediated FCNC is not present at tree level [139].

Note that we also introduce a new discrete symmetry Z_2^{ex} in table 1, and assign positive parity to the SM fermions and the negative parity to extra fermions. Φ and $H_{1,2}$ are Z_2^{ex}-even. As discussed in ref. [140], the Z_2^{ex} symmetry is defined by the subgroup of U(1)$_\chi$ or U(1)$_\psi$ multiplied by $(-1)^{2s}$, where s is the spin of the matter fields.[1] That is, it is anomaly-free and can be realized by one proper Higgs field to break the GUT symmetry. From the phenomenological point of view, this Z_2^{ex} symmetry is required to avoid the tree-level FCNCs, forbidding the tree-level mass mixing between the SM quarks and the extra quarks.

Below, we explain the relevant couplings in our U(1)$'$ model, where there are only two Higgs doublets and Φ, together with the chiral fermions for the anomaly-free conditions, as shown in table 1.

2.2 Extra leptophobic gauge boson

We consider a U(1)$'$ model with two Higgs doublets, where only H_1 carries nonzero U(1)$'$ charge. The kinetic term of H_1 has an extra term associated with the U(1)$'$ gauge boson $\hat{Z}_H$,

$$D^\mu H_1 = D_\mu^{\text{SM}} H_1 - i g_H \hat{Z}_H^\mu H_1 \tag{2.1}$$

with being the U(1)$'$ gauge coupling g_H, while that of H_2 has only the SM part. The mass matrix of $\hat{Z}$ and $\hat{Z}_H$ is given by

$$\mathcal{M}_{\hat{Z},\hat{Z}_H}^2 = \begin{pmatrix} g_Z^2 v_H^2 & -g_Z g_H v_1^2 \\ -g_Z g_H v_1^2 & g_H^2 (v_1^2 + v_\Phi^2) \end{pmatrix}, \tag{2.2}$$

where the mixing angle ξ between $\hat{Z}$ and $\hat{Z}_H$ is given by

$$\tan 2\xi = -\frac{2 g_Z g_H v_1^2}{g_H^2 (v_1^2 + v_\Phi^2) - g_Z^2 v_H^2}. \tag{2.3}$$

g_Z is the gauge coupling in the SM: $g_Z = \sqrt{g'^2 + g^2}$, where g' and g are the U(1)$_Y$ and SU(2)$_L$ gauge couplings. The new gauge boson Z_H is constrained by the collider searches and the electroweak precision tests. In ref. [140], the present authors studied them, finding out that $g_H \lesssim 0.1$ for $M_{Z_H} \gtrsim 400\,\text{GeV}$, and $g_H \approx 0.01$ for $M_{Z_H} \sim 200\,\text{GeV}$ (see section IV.C.1 and figure 1 in ref. [140] for more detail). For such a small g_H, the $Z_H Z_H$ fusion into h_Φ would be small, and shall be ignored in this paper.

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[1]This corresponds to the R-parity in the minimal supersymmetric SM.

2.3 Scalar sector

In our model, there are two Higgs doublets, H_1 and H_2, where H_1 (H_2) gives masses up-type quarks (down-type quarks and charged leptons) after EW symmetry breaking, and one singlet scalar Φ with nonzero U(1)$'$ charge. H_1 is charged under the extra U(1)$'$ symmetry, which could be the origin of the Z_2 symmetry of the usual 2HDM, while H_2 is uncharged. Then, the so-called μ term of the Higgs potential, $\mu H_1^\dagger H_2$ which breaks Z_2 softly, is not invariant under the U(1)$'$ symmetry. This μ-term is replaced by $\mu_\Phi H_1^\dagger H_2 \Phi$, and the μ term will be generated after U(1)$'$ symmetry breaking. Note that 2HDMs with softly broken Z_2 symmetry are discussed widely. This U(1)$'$, which distinguishes two Higgs doublets, suggest the origin of Z_2 symmetry [139] (see ref. [144] for implementation to SU(2)$_H$ gauge symmetry).

The potential of the scalar fields in our model is given by

$$V_{\text{scalar}} = \tilde{m}_1^2 H_1^\dagger H_1 + \tilde{m}_2^2 H_2^\dagger H_2 + \frac{\lambda_1}{2}\left(H_1^\dagger H_1\right)^2 + \frac{\lambda_2}{2}\left(H_2^\dagger H_2\right)^2$$
$$+ \lambda_3 H_1^\dagger H_1 H_2^\dagger H_2 + \lambda_4 H_1^\dagger H_2 H_2^\dagger H_1 + V_\Phi, \tag{2.4}$$

where the potential containing the singlet scalar Φ is

$$V_\Phi = \tilde{m}_\Phi^2 \Phi^\dagger \Phi + \frac{\lambda_\Phi}{2}\left(\Phi^\dagger \Phi\right)^2 - \left(\mu_\Phi H_1^\dagger H_2 \Phi + \text{h.c.}\right) + \tilde{\lambda}_1 H_1^\dagger H_1 \Phi^\dagger \Phi + \tilde{\lambda}_2 H_2^\dagger H_2 \Phi^\dagger \Phi. \tag{2.5}$$

After spontaneous symmetry breaking, the scalar fields can be expanded around their vacuum expectation values as

$$H_i = \begin{pmatrix} \phi_i^+ \\ \frac{1}{\sqrt{2}}(v_i + h_i + i\chi_i) \end{pmatrix}, \quad \Phi = \frac{1}{\sqrt{2}}(v_\Phi + h_\Phi + i\chi_\Phi), \tag{2.6}$$

where $v_1 = v_H \cos\beta$, $v_2 = v_H \sin\beta$ and $v_H = 246\,\text{GeV}$. The neutral CP-even scalars generally mix with each other. The mass matrix is given by

$$\tilde{\mathcal{M}}^2 = \begin{pmatrix} \tilde{\mathcal{M}}_{11}^2 & \tilde{\mathcal{M}}_{12}^2 & \tilde{\mathcal{M}}_{1\Phi}^2 \\ \tilde{\mathcal{M}}_{12}^2 & \tilde{\mathcal{M}}_{22}^2 & \tilde{\mathcal{M}}_{2\Phi}^2 \\ \tilde{\mathcal{M}}_{1\Phi}^2 & \tilde{\mathcal{M}}_{2\Phi}^2 & \tilde{\mathcal{M}}_{\Phi\Phi}^2 \end{pmatrix} \tag{2.7}$$

where

$$\tilde{\mathcal{M}}_{11}^2 = \frac{1}{2}\lambda_1 v_H^2 \cos^2\beta + \frac{1}{2\sqrt{2}}\mu_\Phi v_\Phi \tan\beta, \tag{2.8}$$

$$\tilde{\mathcal{M}}_{22}^2 = \frac{1}{2}\lambda_2 v_H^2 \sin^2\beta + \frac{1}{2\sqrt{2}}\mu_\Phi v_\Phi \cot\beta, \tag{2.9}$$

$$\tilde{\mathcal{M}}_{\Phi\Phi}^2 = \frac{1}{2}\lambda_\Phi v_\Phi^2 + \frac{1}{2\sqrt{2}}\frac{\mu_\Phi v_H^2}{v_\Phi}\sin\beta\cos\beta, \tag{2.10}$$

$$\tilde{\mathcal{M}}_{12}^2 = -\frac{1}{\sqrt{2}}\mu_\Phi v_\Phi + (\lambda_3 + \lambda_4)v_H^2 \sin\beta\cos\beta, \tag{2.11}$$

$$\tilde{\mathcal{M}}_{1\Phi}^2 = \tilde{\lambda}_1 v_H v_\Phi \cos\beta - \frac{1}{\sqrt{2}}\mu_\Phi v_H \sin\beta, \tag{2.12}$$

$$\tilde{\mathcal{M}}_{2\Phi}^2 = \tilde{\lambda}_2 v_H v_\Phi \sin\beta - \frac{1}{\sqrt{2}}\mu_\Phi v_H \cos\beta. \tag{2.13}$$

Since the recent data at the LHC imply that the $125\,\mathrm{GeV}$ scalar boson is almost the SM-like Higgs boson, the mixing between $h_{1,2}$ and h_Φ must be small. For simplicity we assume that there is no mixing between them by setting

$$\tilde{\lambda}_1 = \frac{\mu_\Phi}{\sqrt{2}v_\Phi}\tan\beta, \ \tilde{\lambda}_2 = \frac{\mu_\Phi}{\sqrt{2}v_\Phi}\cot\beta. \tag{2.14}$$

The mass of h_Φ is determined by $m_{h_\Phi}^2 = \tilde{\mathcal{M}}_{\Phi\Phi}^2$.

The other two CP-even Higgs bosons, h_1 and h_2 mix with each other and we identify the light boson as the SM-like Higgs boson h with the mass $m_h = 125\,\mathrm{GeV}$ while the other one is the heavy Higgs boson H:

$$\begin{pmatrix} h_1 \\ h_2 \end{pmatrix} = \begin{pmatrix} \cos\alpha_h & -\sin\alpha_h \\ \sin\alpha_h & \cos\alpha_h \end{pmatrix} \begin{pmatrix} H \\ h \end{pmatrix}, \tag{2.15}$$

with the mixing angle α_h. Since the properties of the Higgs boson measured at the LHC is consistent with the SM prediction, we assume that the decoupling limit, $\sin(\beta - \alpha_h) = 1$, is realized in the two Higgs doublet sector. Note that we still have to keep two Higgs doublets in order to write down the Yukawa couplings for all the observed SM chiral fermions.

2.4 Yukawa sector of extra fermions

Finally, we introduce our Yukawa couplings involving the scalars and the fermions in table 1. The U(1)$'$-symmetric Yukawa couplings in our model are given by

$$V_y = y_{ij}^u \overline{u_R^j} H_1^\dagger i\sigma_2 Q^i + y_{ij}^d \overline{d_R^j} H_2 Q^i + y_{ij}^e \overline{e_R^j} H_2 L^i + y_{ij}^n \overline{n_R^j} H_1^\dagger i\sigma_2 L^i + \mathrm{H.c.}, \tag{2.16}$$

where σ_2 is the Pauli matrix. The Yukawa couplings to generate the mass terms for the extra particles are

$$V^{\mathrm{ex}} = y_{ij}^D \overline{D_R^j} \Phi D_L^i + y_{ij}^H \overline{\widetilde{H}_R^j} \Phi \widetilde{H}_L^i + y_{IJ}^N \overline{N_L^c} H_1^\dagger i\sigma_2 \widetilde{H}_L^i + y_{IJ}'^N \overline{\widetilde{H}_R^i} H_2 N_L^j + \mathrm{H.c.} \ . \tag{2.17}$$

Let us comment on the mass spectrum derived from V_y and V^{ex}. H_1, H_2 and Φ develop nonzero VEVs, and break $\mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y$ and $\mathrm{U}(1)'$ symmetries. The extra colored and charged particles obtain heavy masses from the nonzero VEV of Φ. We also find the neutral particle masses are generated by the VEVs of Higgs doublets and Φ. These massive extra particles are Z_2^{ex}-odd, and thus the lightest neutral fermion among them becomes stable and could be a good cold dark matter candidate [140]. The detailed phenomenological study of the fermionic DM ψ_X (which is mostly n_L) scenario is presented in ref. [140]. As we mentioned in subsection 2.1 and ref. [140], Z_2^{ex} is the subgroup of E_6, which is either $\mathrm{U}(1)_\psi$ or $\mathrm{U}(1)_\chi$, and might also be generated by the E_6 gauge group.

On the other hand, the charged extra leptons decay to the extra neutral particles and charged leptons, and the colored extra quarks decay to the extra neutral ones and the SM particles through the higher-dimensional operators [140]. The direct search for the extra particles at the LHC imposes the lower bounds on their masses. Their signals are colored or charged particles with large missing energy, so that the current lower mass bounds are about $400-800\,\mathrm{GeV}$ [145–147]. However we have to keep in mind that these bounds depend on the dark matter mass, and thus are quite model dependent.

2.5 Scalar DM

Before moving to our phenomenology, let us suggest another possibility of the dark matter, adding an extra scalar, X. Without any theoretical problem, one can introduce a new Z_2^{ex}-odd scalar field X with the $\text{SU}(3)_C \times \text{SU}(2)_L \times \text{U}(1)_Y \times \text{U}(1)'$ quantum numbers equal to $(1, 1, 0; -1)$. The quantum number of X is the same as $\overline{n_R^i}$, and the Z_2^{ex} charge of X is given by the Z_2^{ex} charge of the neutrino multiplied by $(-1)^{2s}$. That is, X could be interpreted as the superpartner of $\overline{n_R^i}$, in the framework of the supersymmetric E_6 GUT.

The gauge-invariant Lagrangian involving X is given by

$$
\begin{aligned}
\mathcal{L}_X = {}& D_\mu X^\dagger D^\mu X - (m_{X0}^2 + \lambda_{H_1 X} H_1^\dagger H_1 + \lambda_{H_2 X} H_2^\dagger H_2) X^\dagger X - \lambda_X (X^\dagger X)^2 \\
& - \left(\lambda_{\Phi X}'' (\Phi^\dagger X)^2 + \text{H.c.} \right) - \lambda_{\Phi X} \Phi^\dagger \Phi X^\dagger X - \lambda_{\Phi X}' |\Phi^\dagger X|^2 \\
& - \left(y_{dX}^D \overline{d_R} D_L X + y_{LX}^{\tilde{H}} \overline{L} \widetilde{H}_R X^\dagger + \text{H.c.} \right).
\end{aligned}
\tag{2.18}
$$

Generation indices are suppressed for simplicity, but should be included in actual calculation. We have imposed Z_2^{ex} symmetry, which forbids dangerous terms such as

$$
\Phi^\dagger X, \ \ H_1^\dagger H_1 \Phi^\dagger X, \ \text{etc.}
$$

that would make X decay. Assuming that $\langle X \rangle = 0$, X would be stable and make another good candidate of CDM, in addition to a neutral fermion DM discussed in the previous subsection. Since X feels $\text{U}(1)'$ gauge force, it is a baryonic DM and interact with the nuclei through Z_H exchanges.

The exotic quark D_L will decay into $\bar{d}_R + \bar{X}$ through the y_{dX}^D term. The collider signature will be dijet + missing E_T and is similar to the squark search bounds. Likewise, the exotic leptons $\tilde{H}_R$ can decay into $l + X$. The collider signature will be dilepton + missing E_T and is similar to the slepton search bounds. Note that the fermionic DM ψ_X (mostly composed of n_L) discussed in the previous subsection can decay into $X + \nu$ if it is kinematically allowed. Therefore the lighter one of X and $\psi_X \approx n_L$ would be a good DM candidate.

3 Phenomenology

As we have mentioned, we assume that the diphoton excess around $750\,\text{GeV}$ is interpreted as the resonant production of the CP-even scalar h_Φ. In the simple setup in section 2, h_Φ does not interact with the SM fermions at the tree level since there is no mixing with h_i. The ATLAS data prefer rather a large decay width ($\Gamma \sim 45\,\text{GeV}$) of h_Φ, but we note that the narrow width of h_Φ where $\Gamma_\Phi \lesssim 1\% \times m_\Phi$ is not excluded either by the data. In fact the CMS data favor a narrower width, in contrast to the ATLAS data. This issue is very important for the model buildings when the $750\,\text{GeV}$ diphoton excess is confirmed in the future.

In the following analysis, we will simply assume that $\Gamma_{\text{tot}} = 10\,\text{GeV}$. We note that the decay width for $h_\Phi \to gg$, $\gamma\gamma$, or any SM particles is much less than $O(10)\,\text{GeV}$. In order to enhance the decay width, the exotic decay channels of h_Φ will be assumed. One

of the candidate is the h_Φ decay into a pair of fermion DM, which exists in the model naturally. The decay into a pair of the scalar DM, which is introduced in section 2.5, could be another candidate.

When the gluon fusion is dominant for the h_Φ production, the cross section for the diphoton production via the h_Φ resonance can be described in terms of the decay widths of $h_\Phi \to gg$ and $h_\Phi \to \gamma\gamma$ and the integral of parton distribution functions (pdfs) of gluons (C_{gg}) by

$$\sigma(gg \to h_\Phi \to \gamma\gamma) = \frac{C_{gg}}{sm_{h_\Phi}\Gamma_{\text{tot}}}\Gamma[h_\Phi \to gg]\Gamma[h_\Phi \to \gamma\gamma], \tag{3.1}$$

where

$$C_{gg} = \frac{\pi^2}{8}\int_\tau^1 \frac{dx}{x}g\left(x, m_\Phi^2\right)g\left(\frac{\tau}{x}, m_\Phi^2\right), \tag{3.2}$$

with $\tau = m_\Phi^2/s$ and $g(x, Q^2)$ is the gluon pdf at $x = Q^2$. Numerically, $C_{gg} = 2137$ for LHC@13TeV and $C_{gg} = 174$ for LHC@8TeV [23].

The decay rates of h_Φ to two gluons and two photons are given by

$$\Gamma[h_\Phi \to gg] = \frac{\alpha_s^2 m_{h_\Phi}^3}{128\pi^3 v_\Phi^2}\left|\sum_{q'}A_{1/2}^H(\tau_{q'})\right|^2, \tag{3.3}$$

$$\Gamma[h_\Phi \to \gamma\gamma] = \frac{\alpha^2 m_{h_\Phi}^3}{256\pi^3 v_\Phi^2}\left|\sum_{q'}N_cQ_{q'}^2 A_{1/2}^H(\tau_{q'}) + \sum_{l'}Q_{l'}^2 A_{1/2}^H(\tau_{l'}) + \frac{v_H v_\Phi}{2m_{H^\pm}^2}\lambda_{h_\Phi H^+ H^-}A_0^H(\tau_{H^\pm})\right|^2, \tag{3.4}$$

where $\tau_i = m_{h_\Phi}^2/4m_i^2$ and $q'(l')$ are the extra quarks (charged leptons), respectively. The loop functions are defined by

$$A_{1/2}^H(\tau) = 2[\tau + (\tau - 1)f(\tau)]/\tau, \tag{3.5}$$

$$A_0^H(\tau) = -[\tau - f(\tau)]/\tau^2, \tag{3.6}$$

where the fucntion $f(x)$ is defined by

$$f(x) = \begin{cases} \arcsin^2\sqrt{x} & , \text{ for } x \leq 1; \\ -\frac{1}{4}\left[\log\frac{1 + \sqrt{1 - 1/x}}{1 - \sqrt{1 - 1/x}} - i\pi\right]^2 & , \text{ for } x > 1. \end{cases} \tag{3.7}$$

We note that there is no W-loop contribution to $h_\Phi \to \gamma\gamma$ since h_ϕ does not couple with the W boson at the tree level. However the charged Higgs boson contributes to the two-photon decay width via the charged Higgs loop with the $h_\Phi H^+ H^-$ coupling

$$\lambda_{h_\Phi H^+ H^-} = \tilde{\lambda}_1 \sin^2\beta + \tilde{\lambda}_2 \cos^2\beta + \sqrt{2}\mu_\Phi \sin\beta\cos\beta/v_\Phi, \tag{3.8}$$

normalized by v_Φ. The charged Higgs contribution to $h_\Phi \to \gamma\gamma$ gets smaller when the charged Higgs boson becomes heavier, showing the typical decoupling behavior.

We find that in a reasonable parameter set, the charged Higgs contribution is not so large. However, its contribution could become more important if the extra fermions get larger masses. In this work, we do not consider the charged Higgs contribution, which is more model-dependent.

Numerically, the decay width of $h_\Phi \to gg$ could be $\mathcal{O}(10)\,\mathrm{GeV}$ for large Yukawa coupling $Y \approx 5 - 10$ and small m_f (exotic fermion mass). But for $y = 1$, it is less than about $0.5\,\mathrm{GeV}$. On the other hand, the decay width of $h_\Phi \to \gamma\gamma$ is at most $\mathcal{O}(0.1)\,\mathrm{GeV}$ even for large Yukawa coupling and small m_f. For $y = 1$, the decay width is less than $5 \times 10^{-3}\,\mathrm{GeV}$ in the entire region of m_f. Without extra decay channels of h_Φ, we cannot achieve $O(10)\,\mathrm{GeV}$ decay width for h_Φ for large m_f ($\gtrsim 500\,\mathrm{GeV}$).

The decay rate for $h_\Phi \to Z\gamma$ is given by

$$\Gamma[h_\Phi \to Z\gamma] = \frac{\alpha m_W^2 m_{h_\Phi}^3}{128\pi^2 v_H^2 v_\Phi^2} \left(1 - \frac{m_Z^2}{m_{h_\Phi}^2}\right)^3 \left| \sum_{f'} N_{f'} \frac{Q_{f'} v_{f'}}{c_w} A_{1/2}^H(\tau_{f'}, \lambda_{f'}) \right|^2, \qquad (3.9)$$

where $N_{q'} = N_c$, $N_{l'} = 1$, and $v_f = 2I_f^3 - 4Q_f x_w$. Here we do not consider the contribution of the charged Higgs boson like the $h_\Phi \to \gamma\gamma$ decay.

The function $A_{1/2}^H(\tau, \lambda)$ is defined by

$$A_{1/2}^H(\tau, \lambda) = I_1(\tau^{-1}, \lambda^{-1}) - I_2(\tau^{-1}, \lambda^{-1}), \qquad (3.10)$$

where

$$I_1(x, y) = \frac{xy}{2(x - y)} + \frac{x^2 y^2}{2(x - y)^2}\left[f(x^{-1}) - f(y^{-1})\right] + \frac{x^2 y}{(x - y)^2}\left[g(x^{-1}) - g(y^{-1})\right], \quad (3.11)$$

$$I_2(x, y) = -\frac{xy}{2(x - y)}\left[f(x^{-1}) - f(y^{-1})\right] \qquad (3.12)$$

with

$$g(x) = \begin{cases} \sqrt{x^{-1} - 1}\,\arcsin\sqrt{x} & , \quad \text{for } x \leq 1, \\ \frac{\sqrt{1 - x^{-1}}}{2}\left[\log\frac{1 + \sqrt{1 - 1/x}}{1 - \sqrt{1 - 1/x}} - i\pi\right]^2 & , \quad \text{for } x > 1, \end{cases} \qquad (3.13)$$

and $f(x)$ is defined in eq. (3.5).

Besides, h_Φ can decay to the extra particles, as well as the dark matter particles. Now, let us simply define the extra decay width $\Delta\Gamma$ (GeV) and the total decay width (Γ_{tot}) of h_Φ could be given by

$$\Gamma_{\mathrm{tot}} = \Delta\Gamma + \Gamma[h_\Phi \to gg] + \Gamma[h_\Phi \to \gamma\gamma] + \Gamma[h_\Phi \to \gamma Z]. \qquad (3.14)$$

The diphoton excess requires $\mathcal{O}(10)$-GeV Γ_{tot} and $\mathcal{O}(10)$ fb diphoton signal at $\sqrt{s} = 13\,\mathrm{TeV}$. This means that large $\Delta\Gamma$ is necessary to reproduce the excess. In figure 1, we see the required Yukawa coupling ($y = \sqrt{2}m_f/v_\Phi$), where m_f is the mass of the extra fermions, and the diphoton signal at LHC13. In the left panel, the total decay width is fixed at $10\,\mathrm{GeV}$, which can be readily achieved by allowing the invisible decay of h_Φ into a pair of DM particles (see figure 2 and the related discussions). In the right, $\Gamma_{\mathrm{tot}} = 1\,\mathrm{GeV}$ (pink) and $\Gamma_{\mathrm{tot}} = 10\,\mathrm{GeV}$ (cyan) are shown, when m_f is between $500\,\mathrm{GeV}$ and $1\,\mathrm{TeV}$. Note that we need large Yukawa coupling $y \approx 5 - 10$ for $m_f > 400\,\mathrm{GeV}$ in order to get the correct size of the production cross section for $pp \to h_\Phi \to \gamma\gamma$. Even if the total decay width is $1\,\mathrm{GeV}$, we still need large Yukawa coupling, as we see in the right panel.
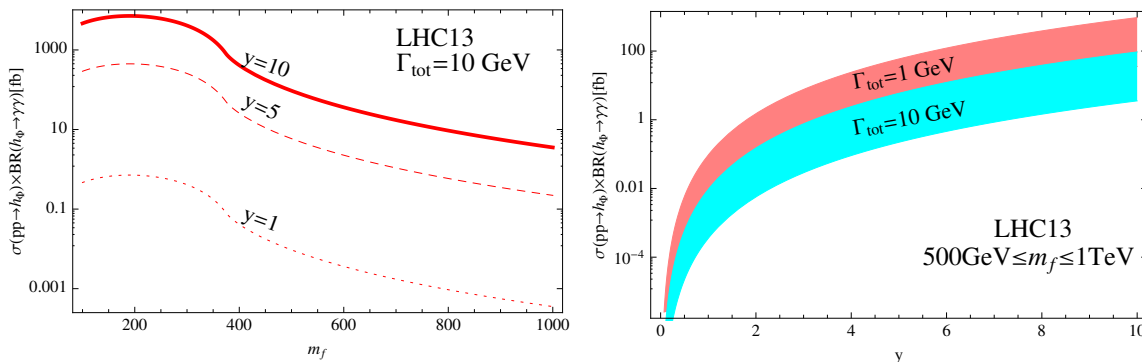
Figure 1. (Left): m_f vs. the diphoton signal via the gluon-gluon fusion at the LHC@13TeV. The total decay widths is fixed at $\Gamma_{\rm tot} = 10\,{\rm GeV}$. (Right): y vs. the diphoton signal for different values of the $\Gamma_{\rm tot}$ and for $500\,{\rm GeV} \leq m_f \leq 1\,{\rm TeV}$.

Now, we discuss the extra decay width $\Delta\Gamma$ in greater detail. First of all, h_Φ can decay into $Z_H Z_H$, a pair of leptophobic U(1)$'$ gauge bosons. If there exists the mixing between $\hat{Z}$ and $\hat{Z}_H$ bosons, it can also decay into ZZ and ZZ_H through the mixing at the tree level. But the decay widths will be suppressed by small gauge coupling $g_H \lesssim O(0.1)$ and the small $Z - Z_H$ mixing, which are required from the search for the extra neutral gauge boson [140]. There is also contribution to $h_\Phi \to WW, ZZ, ZZ_H, Z_H Z_H$ from the loop diagrams of extra fermions which have SU(2)$_L$ and/or U(1)$'$ quantum numbers. The last two channels are suppressed by the small g_H coupling, while the decay widths in the first two channels may be the same order as $\Gamma[h_\Phi \to \gamma\gamma]$. Numerically, their contribution to the total decay width would be negligible like $\Gamma[h_\Phi \to \gamma\gamma]$ in the case of $\Gamma_{\rm tot} \sim O(10)\,{\rm GeV}$.

Next, we consider the h_Φ decay into two scalar bosons $(s_i s_j)$, $s_i = h, H, H^\pm, A$, and into two Majorana fermions $(\psi_X \psi_X)$, where ψ_X is one of dark matter candidates, if kinematically allowed. These new channels are unique aspects of our model due to the 2HDM nature of leptophobic U(1)$'$ models. We can enhance the decay width of h_Φ not by ad hoc way but by underlying gauge theory which is anomaly-free. The decay widths for $h_\Phi \to s_i s_j$ and $h_\Phi \to \psi_X \psi_X$ are given by

$$\Gamma[h_\Phi \to s_i s_i] = \frac{\lambda_{h_\Phi s_i s_i}^2}{32\pi} \frac{v_\Phi^2}{m_{h_\Phi}} \sqrt{1 - 4x_i}, \tag{3.15}$$

$$\Gamma[h_\Phi \to s_i s_j] = \frac{\lambda_{h_\Phi s_i s_j}^2}{16\pi} \frac{v_\Phi^2}{m_{h_\Phi}} \lambda^{1/2}(1, x_i, x_j) (\text{for } i \neq j) , \tag{3.16}$$

$$\Gamma[h_\Phi \to \psi_X \psi_X] = \frac{y_X^2 m_{h_\Phi}}{8\pi} \{1 - 4(m_{\psi_X}^2/m_{h_\Phi}^2)\}^{\frac{3}{2}}, \tag{3.17}$$

where $x_i = m_{s_i}^2/m_{h_\Phi}^2$ and $\lambda(a, b, c) = a^2 + b^2 + c^2 - 2ab - 2bc - 2ca$. y_X and m_{ψ_X} are the Yukawa coupling with h_Φ and the mass of ψ_X.

The non-zero entries of $\lambda_{h_\Phi s_i s_j}$ are given by

$$\lambda_{h_\Phi hh} = \tilde{\lambda}_1 \sin^2 \alpha + \tilde{\lambda}_2 \cos^2 \alpha + \frac{\mu_\Phi}{\sqrt{2} v_\Phi} \sin \alpha \cos \alpha, \tag{3.18}$$

$$\lambda_{h_\Phi HH} = \tilde{\lambda}_1 \cos^2 \alpha + \tilde{\lambda}_2 \sin^2 \alpha - \frac{\mu_\Phi}{\sqrt{2} v_\Phi} \sin \alpha \cos \alpha, \tag{3.19}$$

$$\lambda_{h_\Phi hH} = -\tilde{\lambda}_1 \sin 2\alpha + \tilde{\lambda}_2 \cos 2\alpha - \frac{\mu_\Phi}{\sqrt{2}v_\Phi} \cos 2\alpha, \tag{3.20}$$

$$\lambda_{h_\Phi AA} = \frac{v_\Phi^2}{v_\Phi^2 + (v_H \sin\beta \cos\beta)^2}$$
$$\times \left(\tilde{\lambda}_1 \sin^2\beta + \tilde{\lambda}_2 \cos^2\beta + \frac{\mu_\Phi}{\sqrt{2}v_\Phi} \sin\beta \cos\beta + 2\lambda_\Phi v_H \sin\beta \cos\beta \right), \tag{3.21}$$

where A is the pseudoscalar boson. The coupling $\lambda_{h_\Phi H^+ H^-}$ has been defined in eq. (3.5) in the context of $h_\Phi \to \gamma\gamma$.

In addition, we find a dark matter candidate among the natural fermions [140]. Assuming that the Yukawa couplings are flavor-independent, we can explicitly calculate the DM mass and the Yukawa coupling with h_Φ. Figure 2 shows the partial decay widths of h_Φ to two dark matter particles in the fermionic DM scenario (left) and scalar DM scenario (right). As we see, the h_Φ invisible decay cannot be so large in the fermionic DM case, although the branching ratio is relatively larger than the diphoton decay width. In the region where the perturbativity holds, the invisible decay width of h_Φ is at most about $10\,\mathrm{GeV}$. In this case, if the total decay width of the $750\,\mathrm{GeV}$ excess is confirmed to be about $45\,\mathrm{GeV}$, then other decay channels like $h_\Phi \to hH, HH, AA$ must be comparable to or dominant over the invisible decay[2] (see the end of this section for more discussion on this point). On the other hand, the invisible decay in the scalar DM case can be dominant, if $\lambda_{\Phi X}$ is $\mathcal{O}(1)$.

In our model, there is a massive gauge boson (Z_H), which dominantly gets the mass from the nonzero VEV of Φ. The extra fermions also get the mass from the VEV, so there is a relation between the Z_H mass and m_f. Approximately, it can be evaluated as

$$M_{Z_H} \approx \frac{g_H}{y} m_f. \tag{3.22}$$

The diphoton excess suggests $\mathcal{O}(1)$ y, and then M_{Z_H} is at least $\mathcal{O}(100)\,\mathrm{GeV}$, because g_H could not be $\mathcal{O}(1)$ to evade the stringent bound from the dijet signal. Another strong constraint is from ρ parameter, as discussed in ref. [140]. In this scenario, Z_H is light, so the coupling should be small. As discussed in refs. [139–141], the ρ parameter is deviated from 1 at the tree-level, because of Z-Z_H mixing. The bound is roughly estimated as [141]

$$\frac{g_H}{g_Z} \frac{M_Z^2}{|M_{Z_H}^2 - M_Z^2|} \lesssim 0.004. \tag{3.23}$$

Then $\mathcal{O}(100)$-GeV $M_{Z_H}^2$ requires $g_H \lesssim \mathcal{O}(0.1) \times g_Z$, which may be too small to enhance the branching ratio of $h_\Phi \to Z_H Z_H$.

There are some experimental constraints relevant to our scenario. Since the h_Φ is produced from the gluon fusion copiously at the LHC, the dijet production can severely constrain our scenario. The bound on the dijet production at LHC@8TeV is about 2 pb at CMS [149]. By imposing

$$\sigma(gg \to h_\Phi \to gg) \lesssim 2 \text{ pb}, \tag{3.24}$$

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[2]The charged scalar mass is constrained by $B \to X_s\gamma$ in type-II Higgs doublet model, and should satisfy $m_{H^\pm} \gtrsim 500\,\mathrm{GeV}$ [148]. Therefore $h_\Phi \to H^+ H^-$ is kinematically forbidden in our model.

we find that the mass of the exotic quarks should be larger than $400\,\mathrm{GeV}$ for the Yukawa coupling $y = 5$ and $600\,\mathrm{GeV}$ for $y = 10$.

Next, we consider the diboson channels. First, the h_Φ can decay into hh. Then this channel is constrained by the experimental data at LHC@8TeV [150]:

$$\sigma(gg \to h_\Phi \to hh) \lesssim 10 \ \mathrm{fb}. \tag{3.25}$$

The decay width of $h_\Phi \to hh$ strongly depends on the model parameters. For example, for $\alpha_h = 0$ and $\mu_\Phi \sim v_H$, the branching ratio of $h_\Phi \to hh$ is $\mathcal{O}(0.1)$ for $\tan\beta \sim 1$, while it could be $\mathcal{O}(0.01)$ for $\tan\beta \sim 10$. Actually, the bound (3.25) requires $Br(h_\Phi \to hh) \lesssim 0.01$.

We note that there are other diboson channels, $h_\Phi \to ZZ, WW$. The bounds from LHC8 for the WW production [151] and ZZ production [152] are

$$\sigma(gg \to h_\Phi \to WW) \lesssim 40 \ \mathrm{fb}, \tag{3.26}$$

$$\sigma(gg \to h_\Phi \to ZZ) \lesssim 10 \ \mathrm{fb}, \tag{3.27}$$

respectively. In our model, h_Φ does not interact with W and Z bosons at the tree-level directly so that the decay channels are suppressed by loop diagrams or Z-Z_H mixing, which must be small. Since the loop diagrams of the extra leptons which has the SM $SU(2)_L$ quantum number are dominant, the decay width for $h_\Phi \to WW$ and ZZ would be the same order as $\gamma\gamma$ or less. Therefore, the bound for the diboson channels WW and ZZ would be acceptable.

Finally, we consider the $Z\gamma$ production channel. At LHC8, the bound on the production is [153]

$$\sigma(gg \to h_\Phi \to Z\gamma) \lesssim 3.8 \ \mathrm{fb}. \tag{3.28}$$

From the numerical analysis, we find that the extra fermion mass should be larger than about 200 (400) GeV for $y = 5\,(10)$, respectively. For $y = 1$, the cross section is less than $1\,\mathrm{fb}$. They are weaker than the bound from the dijet production.

If the invisible decay of h_Φ is dominant, the monojet search at the LHC would give most stringent constraints on the models which may explain the diphoton excess. The NLO correction to $gg \to h_\Phi$ is very involved, and beyond the scope of this paper. Here we try to make a qualitative argument on the monojet $+ \not{E}_T$ constraints. The monojet $+ \not{E}_T$ signal will be generated by the parton level processes: (i) the initial state radiation of gluon in $gg \to h_\Phi$ and $qg \to qgg$ followed by $gg \to h_\Phi$ via triangle diagrams and (ii) $gg \to h_\Phi g$ via box diagrams. The type (i) will mainly generate a monojet in the beam direction with low p_T and may be removed by the $\not{E}_T$ cut. They are also suppressed by an extra g_s. The type (ii) could generate high p_T monojet and should be constrained by the data on monojet $+ \not{E}_T$. If we use a naive dimensional analysis, its rate would be suppressed by $\sim \alpha_s/(4\pi)$ compared with the rate for $gg \to h_\Phi$ from a triangle diagram. At the LHC@8TeV, $\sigma(gg \to h_\Phi) \approx 2$ pb, so that we would expect that $\sigma(gg \to h_\Phi g) \approx 0.2$ pb, which satisfies the bound $\lesssim 0.8$ pb derived in ref. [23]. For more definitive conclusion on this issue, we have to perform more detailed analysis.
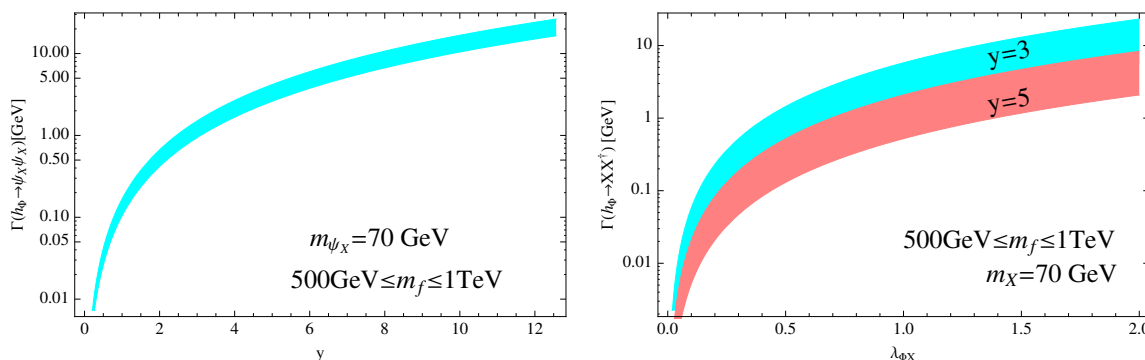
Figure 2. y vs. invisible decay width of h_Φ (GeV) in the fermionic DM scenario (left) and scalar DM scenario (right).The vector-like fermion mass is between $500\,\mathrm{GeV}$ and $1\,\mathrm{TeV}$ on the cyan and pink bands. The dark matter masses are $70\,\mathrm{GeV}$ in the both cases.

The dijet + $\not{E}_T$ process may constrain our model. For the vector boson fusion process, this dijet + $\not{E}_T$ occurs via a parton level process,

$$q\bar{q}' \to q\bar{q}' + Z_H Z_H, \quad \text{followed by}$$
$$Z_H Z_H \to h_\Phi \to XX^\dagger,$$

where X is a dark matter particle. Considering the current bounds on the m_{Z_H} and g_H, this $Z_H Z_H$ fusion production could be neglected safely.

Another dijet + $\not{E}_T$ events could arise from gg fusion to h_Φ beyond the leading order,

$$gg \to gg + gg(\to h_\Phi)$$
$$qq(\text{or } q\bar{q}) \to qq(\text{or } q\bar{q}) + gg(\to h_\Phi)$$
$$qg(\text{or } \bar{q}g) \to qg(\text{or } \bar{q}g) + gg(\to h_\Phi)$$

which will be $O(\alpha_s)$ suppressed compared with what we have studied in the earlier part of this section, namely $gg \to h_\Phi$. And the dijets in these processes will be mainly in the beam directions with low p_T and we expect that they will be removed by the p_T cuts.

Finally, let us comment on other decay modes of h_Φ to two extra scalars, such as $h_\Phi \to Hh$, HH, and AA, whose decay rates eqs. (3.15) and (3.16) are determined by the dimensionless couplings in the Higgs potential, eqs. (3.18)–(3.21). In principle, these extra decay could be sizable up to $\Delta\Gamma \approx \mathcal{O}(10)\,\mathrm{GeV}$. H and A mainly decay into the $b\bar{b}$ state, and so the final states in the $h_\Phi \to Hh$ decay channel would be $bb\bar{b}\bar{b}$. This channel is not strongly constrained by the present data, and is one of the promising signals of our scenario.

4 Summary

In this paper we interpret the recently reported diphoton excess at $750\,\mathrm{GeV}$ in terms of a new singlet scalar boson h_Φ that originates from spontaneous breaking of leptophobic $\mathrm{U}(1)'$ embedded in E_6 grand unification. A **27**-dimensional fundamental representation of E6 gauge group contains one family of SM chiral fermions, as well as 11 more chiral

fermions, some of which are vectorlike under the SM $SU(2)_L \times U(1)_Y$ gauge symmetry. Anomaly cancellation is automatic in this model, and exotic fermions are chiral under $U(1)'$ so that their masses arise entirely from spontaneous breaking of $U(1)'$ symmetry by the nonzero VEV of Φ. The observed diphoton excesses are attributed to $gg \to h_\Phi \to \gamma\gamma$. The vectorlike exotic fermions are chiral under new $U(1)'$ gauge symmetry and their masses are generated only by spontaneous gauge symmetry. Therefore their loop effects would be protected from the decoupling theorem, like the top quark loop contributions to $h \to gg, \gamma\gamma$, etc.. In our model, h_Φ can decay into a pair of DM, as well as two scalar bosons such as hh, Hh, AA, etc.. In particular the Hh final state can have $O(10)\,\mathrm{GeV}$ decay width, making one of the dominant decay channels of h_Φ.

If the diphoton excess at $750\,\mathrm{GeV}$ with large decay width $\sim 45\,\mathrm{GeV}$ is confirmed in the next LHC run, our model predicts that there should be new vectorlike quarks and leptons around $\sim O(1)\mathrm{TeV}$ (or lighter for vectorlike leptons), whose collider signatures would be similar to the squark/slepton searches within the R-parity conservation, namely dijet $+ \not{E}_T$ or dilepton $+ \not{E}_T$. Also additional scalar bosons will be present too, and one of the main decay channels of h_Φ would be Hh final state. In our model the production cross sections for exotic fermions will be larger than the sfermion productions because they are spin-1/2 fermions. In addition, there will be a new leptophobic (baryonic) gauge boson Z_H whose mass could be still as low as a few GeV for a few hundred GeV v_Φ and $g_H \lesssim O(0.1)$, which is mainly constrained by the decay width of the Z boson, Γ_Z and the ρ parameter [140]. DM will be either spin-1/2 fermion or spin-0 scalar, and they will be baryonic in a sense that they have interactions with the nuclei through Z_H exchanges. DM phenomenology within this model in the context of $750\,\mathrm{GeV}$ diphoton excess will be presented elsewhere.

Acknowledgments

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