KIAS-16002

Diphoton Excess at 750 GeV in leptophobic $\mathrm{U}(1)'$ 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 $\mathrm{U}(1)'$ model inspired by the E_6 grand unified theory (GUT). In this model, the Standard Model (SM) chiral fermions carry charges under extra $\mathrm{U}(1)'$ gauge symmetry which is spontaneously broken by a $\mathrm{U}(1)'$-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 $\mathrm{U}(1)'$, 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.

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 Model (BSM). 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]. The local (global) significances are $3.6\,(2.0)\,\sigma$ (ATLAS) and $2.6\,(1.2)\,\sigma$ (CMS), respectively. They are not still conclusive, but it may be important for theorists to survey the possibility of the diphoton resonance in BSMs [2–13].

One of the attractive candidates for the BSMs which predicts the heavy scalar resonance is $\mathrm{U}(1)'$ extension of the Standard Model (SM). In a minimal setup, an additional gauged $\mathrm{U}(1)'$ is introduced and spontaneously broken by the nonzero vacuum expectation value (VEV) of an extra $\mathrm{U}(1)'$-charged scalar (Φ). In the effective lagrangian, an extra massive gauge boson is predicted and a CP-even scalar also appear around the $\mathrm{U}(1)'$ 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, $\mathrm{U}(1)'$ may cause the gauge anomaly depending on the $\mathrm{U}(1)'$ charge assignments to the SM fermions. Only anomaly-free and generation-independent $\mathrm{U}(1)'$ charge assignment is the linear combination of the hypercharge and $\mathrm{U}(1)_{\mathrm{B-L}}$, introducing the right-handed neutrinos. However, such $\mathrm{U}(1)'$ faces a strong bound from the Drell-Yan process, so that the $\mathrm{U}(1)'$ breaking scale can not be low.

As another possibility, we can discuss a $\mathrm{U}(1)'$ model, where the anomaly-free conditions are satisfied by introducing extra chiral fermions. One interesting $\mathrm{U}(1)'$ extension would be the one motivated by the Grand Unified Theory (GUT) with a large rank. For instance, the E_6 GUT principally predict two additional $\mathrm{U}(1)'$ 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 $\mathrm{U}(1)'$. Interestingly, it is suggested that a certain linear combination of the $\mathrm{U}(1)'$ charges could be leptophobic and, in fact, has been widely studied so far in the many literatures [14–20]. Such a leptophobic $\mathrm{U}(1)'$ has been paid an attention to so far, because it can evade the strong bound from the Drell-Yan. Then the $\mathrm{U}(1)'$ breaking scale can be lower than the $\mathrm{U}(1)_{B-L}$ case. In Ref. [18], the present authors discussed the Higgs physics in the leptophobic $\mathrm{U}(1)'$ model. They also found a neutral fermion (ψ_X) dark matter candidate in this model. The 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 Φ, but not Higgs doublets because of the $\mathrm{U}(1)'$ charge assignment. Then, the CP-even scalar mode (h_Φ), which appears after the $\mathrm{U}(1)'$ symmetry breaking, couples only with extra quarks, although it may mix with the other scalars from the 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 Sec. 2, we introduce our model setup inspired by E_6 GUT. In Sec. 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 Sec. 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. [17, 18], where leptophobic U(1)$'$ gauge symmetry is introduced to the SM 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. One simple 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. The E_6 GUT predicts additional two U(1)$'$ and extra chiral fermions: three families of extra right-handed down-type quarks and extra left-handed leptons. Besides, a scalar field, Φ, to break U(1)$'$ is naturally introduced by the **27**-dimensional fields. In this letter, we do not touch the detail of the underling theory, and discuss U(1)$'$ model, where the charge assignment is inspired by the E_6 GUT. Let us introduce the minimal setup of our model below.

In our model, extra U(1)$'$ charges are assigned to quarks but not to left-handed and right-handed charged leptons. Therefore the new U(1)$'$ is leptophobic with respect to the SM fermions. Since this extra symmetry originates from the E$_6$ symmetry, it does not suffer from gauge anomalies, with the introduction of extra chiral fermions which are vector-like under the SM gauge group (see for example [21]).

The quantum numbers of the SM and the extra chiral fermions are summarized in Table 1. Q^i, d^i_R and u^i_R are the SM quarks, which carry nonzero U(1)$'$ charges, while L^i and e^i_R are the SM leptons are not charged under the U(1)$'$, as shown in the table. In addition, extra quarks and leptons are contained in the **27**, which we denote by D^i_L, D^i_R ($i = 1, 2, 3$) and $\widetilde{H}^i_L$, $\widetilde{H}^i_R$ ($i = 1, 2, 3$). We also find the SM-singlet fermions and scalar, n^i_R, N^i_L, and Φ, whose U(1)$'$ charges are -1 *.

There are two Higgs doublets, denoted by H_1 and H_2, which break the electroweak

* Note that all of the fields are embedded into a **27**-representational superfield for the each generation in the supersymmetric extension of this model.

Table 1: Matter contents in U(1)′ model inspired by E_6 GUTs. Here, i denotes the generation index: $i = 1, 2, 3$.

Fields	SU(3)	SU(2)	$\mathrm{U}(1)_Y$	$\mathrm{U}(1)'$	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	0	+
e_R^i	**1**	**1**	−1	0	
n_R^i	**1**	**1**	0	1	
H_2	**1**	**2**	−1/2	0	
H_1	**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	0	−
$\widetilde{H}_R^i$	**1**	**2**	−1/2	−1	
N_L^i	**1**	**1**	0	−1	

symmetry with their nonzero VEVs. Note that there are two Higgs doublets in our model, in order to realize the Yukawa couplings at the renormalizable level, and they carry the different U(1)′ charges to evade the tree-level flavor changing neutral currents. This is an implementation of discrete Z_2 symmetry of the usual 2HDM into continuous $U(1)_H$ gauge symmetry first proposed in Ref. [17].

From the point of view of the top-down approach, one issue may be how to realize the U(1)′ in the low energy regime. The possibility that the leptophobic U(1)′ is generated by kinetic mixing has been studied in Ref. [15]. 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, we will not touch the detail but we simply consider the Z' model inspired by E_6, keeping the minimal set for the anomaly-free conditions.

We have introduced a new discrete symmetry Z_2^{ex} [18], and assigned positive parity to the SM fermions and the negative parity to extra fermions that were introduced for anomaly cancellation.

2.2 Extra leptophobic gauge boson

The kinetic term of H_1 has an extra term associated with the $U(1)_H$ gauge boson $\hat{Z}_H$,

$$D^\mu H_1 = D_\mu^{\rm SM} H_1 - i g_H \hat{Z}_H^\mu H_1 \tag{1}$$

with being the $U(1)_H$ 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

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

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

$$\tan 2\xi = -\frac{2 g_Z g_H v_1^2}{g_H^2 (v_1^2 + v_\Phi^2) - g_Z^2 v_H^2}. \tag{3}$$

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. [18], the present authors studied them, finding out that $g_H \lesssim 0.1$ for $M_{Z_H} \gtrsim 400$ GeV, and $g_H \approx 0.01$ for $M_{Z_H} \sim 200$ GeV (see Sec. IV.C.1 and Fig. 1 in Ref. [18] 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.

2.3 Scalar sector

In our model, there are two Higgs doublets, H_1 and H_2, where H_1 (H_2) give 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 an 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. This way the origin of Z_2 symmetry and its soft breaking in the usual 2HDMs is understood as spontaneous breaking of new Higgs gauge symmetry U(1)$' = U(1)_H = U(1)_b$, which was one of the main motivations for introducing Higgs gauge symmetry in Ref. [17] (see Ref. [22] for implentation to $SU(2)_H$ gauge symmetry).

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

$$\begin{aligned} V_{\rm scalar} = {} & \tilde{m}_1^2 H_1^\dagger H_1 + \tilde{m}_2^2 H_2^\dagger H_2 + \frac{\lambda_1}{2}\left(H_1^\dagger H_1\right)^2 + \frac{\lambda_2}{2}\left(H_2^\dagger H_2\right)^2 \\ & + \lambda_3 H_1^\dagger H_1 H_2^\dagger H_2 + \lambda_4 H_1^\dagger H_2 H_2^\dagger H_1 + V_\Phi, \end{aligned} \tag{4}$$

where the potential contains the singlet scalar Φ is

$$V_\Phi = \tilde{m}_\Phi^2 \Phi^\dagger \Phi + \frac{\lambda_\Phi}{2}\left(\Phi^\dagger\Phi\right)^2 + \left(\mu_\Phi H_1^\dagger H_2 \Phi + {\rm h.c.}\right) + \tilde{\lambda}_1 H_1^\dagger H_1 \Phi^\dagger \Phi + \tilde{\lambda}_2 H_2^\dagger H_2 \Phi^\dagger \Phi. \tag{5}$$

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

$$H_i = \begin{pmatrix} \phi_i^+ \\ \frac{1}{\sqrt{2}}(v_i + h_i + i\chi_i) \end{pmatrix}, \;\; \Phi = \frac{1}{\sqrt{2}}(v_\Phi + h_\Phi + i\chi_\Phi), \tag{6}$$

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

$$\tilde{\mathcal{M}}^2 = \begin{pmatrix} \tilde{\mathcal{M}}^2_{11} & \tilde{\mathcal{M}}^2_{12} & \tilde{\mathcal{M}}^2_{1\Phi} \\ \tilde{\mathcal{M}}^2_{12} & \tilde{\mathcal{M}}^2_{22} & \tilde{\mathcal{M}}^2_{2\Phi} \\ \tilde{\mathcal{M}}^2_{1\Phi} & \tilde{\mathcal{M}}^2_{2\Phi} & \tilde{\mathcal{M}}^2_{\Phi\Phi} \end{pmatrix} \tag{7}$$

where

$$\tilde{\mathcal{M}}^2_{11} = \frac{1}{2}\lambda_1 v_H^2 \cos^2\beta + \frac{1}{2\sqrt{2}}\mu_\Phi v_\Phi \tan\beta, \tag{8}$$

$$\tilde{\mathcal{M}}^2_{22} = \frac{1}{2}\lambda_2 v_H^2 \sin^2\beta + \frac{1}{2\sqrt{2}}\mu_\Phi v_\Phi \cot\beta, \tag{9}$$

$$\tilde{\mathcal{M}}^2_{\Phi\Phi} = \frac{1}{2}\lambda_\Phi v_\Phi^2 + \frac{1}{2\sqrt{2}}\frac{\mu_\Phi v_H}{v_\Phi}\sin\beta\cos\beta, \tag{10}$$

$$\tilde{\mathcal{M}}^2_{12} = -\frac{1}{\sqrt{2}}\mu_\Phi v_\Phi + \lambda_3 v_H^2 \sin\beta\cos\beta + \lambda_4 v_H^2 \cos\beta\sin\beta, \tag{11}$$

$$\tilde{\mathcal{M}}^2_{1\Phi} = \tilde{\lambda}_1 v_H v_\Phi \cos\beta - \frac{1}{\sqrt{2}}\mu_\Phi v_H \sin\beta, \tag{12}$$

$$\tilde{\mathcal{M}}^2_{2\Phi} = \tilde{\lambda}_2 v_H v_\Phi \sin\beta - \frac{1}{\sqrt{2}}\mu_\Phi v_H \cos\beta. \tag{13}$$

Since the recent data at the LHC implies that the 125 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

$$\tilde{\lambda}_1 = \frac{\mu_\Phi}{\sqrt{2}v_\Phi}\tan\beta, \;\; \tilde{\lambda}_2 = \frac{\mu_\Phi}{\sqrt{2}v_\Phi}\cot\beta. \tag{14}$$

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

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 $m_h = 125$ GeV while the other one is the heavy Higgs boson H:

$$\begin{pmatrix} h_1 \\ h_2 \end{pmatrix} = \begin{pmatrix} \cos\alpha_h & -\sin\alpha_h \\ \sin\alpha_h & \cos\alpha_h \end{pmatrix} \begin{pmatrix} H \\ h \end{pmatrix}, \tag{15}$$

with the mixing angle α_h. In the following, we will assume that the observed 750 GeV diphoton excess is mostly composed of h_Φ, and the decoupling limit 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

The U(1)$'$-symmetric Yukawa couplings in our model are given by

$$V_y = y^u_{ij}\overline{u^j_R}H_1^\dagger i\sigma_2 Q^i + y^d_{ij}\overline{d^j_R}H_2 Q^i + y^e_{ij}\overline{e^j_R}H_2 L^i + y^n_{ij}\overline{n^j_R}H_1^\dagger i\sigma_2 L^i + H.c., \tag{16}$$

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

$$V^{\rm ex} = y^D_{ij}\overline{D^j_R}\Phi D^i_L + y^H_{ij}\overline{\widetilde{H}^j_R}\Phi \widetilde{H}^i_L + y^N_{IJ}\overline{N^c_L}H_1^\dagger i\sigma_2 \widetilde{H}^i_L + y'^N_{IJ}\overline{\widetilde{H}^i_R}H_2 N^j_L + H.c. \,. \tag{17}$$

Let us comment on the mass spectrum derived from V_y and $V^{\rm ex}$. H_1, H_2 and Φ develop nonzero VEVs, and break $\mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y$ and 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 $Z_2^{\rm ex}$-odd, and thus the lightest neutral fermion among them becomes stable and could be a good cold dark matter candidate [18]. The detailed phenomenological study of the fermionic DM ψ_X(which is mostly n_L) scenario is presented in Ref. [18]. We note that $Z_2^{\rm ex}$ might be also 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 ones decay to the extra neutral ones and the SM particles through the higher-dimensional operators [18]. 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 $400-800$ GeV [23–25]. 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

One can introduce new $Z_2^{\rm ex}$-odd scalar field X with the $SU(3)_C\times SU(2)_L\times U(1)_Y\times U(1)_H$ quantum numbers equal to $(1,1,0;-1)$. Then the gauge-invariant Lagrangian involving X is given by

$$\begin{aligned}\mathcal{L}_X =\ & D_\mu X^\dagger D^\mu X - (m^2_{X0} + \lambda_{H_1X}H_1^\dagger H_1 + \lambda_{H_2X}H_2^\dagger H_2)X^\dagger X - \lambda_X (X^\dagger X)^2 \\ & - \left(\lambda''_{\Phi X}(\Phi^\dagger X)^2 + H.c.\right) - \lambda_{\Phi X}\Phi^\dagger\Phi X^\dagger X - \lambda'_{\Phi X}|\Phi^\dagger X|^2 \\ & - \left(y^D_{dX}\overline{d_R}D_L X + y^{\tilde{H}}_{LX}\overline{L}\widetilde{H}_R X^\dagger + H.c.\right)\end{aligned} \tag{18}$$

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

$$\Phi^\dagger X\ \ , H_1^\dagger H_1 \Phi^\dagger X,\ \ etc.$$

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 $U(1)_b$ gauge force, it is a baryonic DM and interact with the nuclei through $Z_H = Z_b$ exchanges.

The exotic quark D_L will decay into $\bar{d}_R + \bar{X}$ through the y^D_{dX} 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 GeV is interpreted as the resonant production of the CP-even scalar h_Φ. In the simple setup in Sec. 2, h_Φ does not interact with the SM fermions at the tree level since there is no mixing with h_i.

When the gluon fusion is dominant, 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

$$\sigma(gg \to h_\Phi \to \gamma\gamma) = \frac{C_{gg}}{s m_{h_\Phi} \Gamma_{\rm tot}} \Gamma[h_\Phi \to gg] \Gamma[h_\Phi \to \gamma\gamma], \tag{19}$$

where

$$C_{gg} = \frac{\pi^2}{8} \int_\tau^1 \frac{dx}{x} g\left(x, m_\Phi^2\right) g\left(\frac{\tau}{x}, m_\Phi^2\right), \tag{20}$$

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 [5].

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

$$\Gamma[h_\Phi \to gg] = \frac{\alpha_s^2 m_{h_\Phi}^3}{128\pi^3 v_\Phi^2} \left| \sum_{q'} A^H_{1/2}(\tau_{q'}) \right|^2, \tag{21}$$

$$\Gamma[h_\Phi \to \gamma\gamma] = \frac{\alpha^2 m_{h_\Phi}^3}{256\pi^3 v_\Phi^2} \left| \sum_{q'} N_c Q_{q'}^2 A^H_{1/2}(\tau_{q'}) + \sum_{l'} Q_{l'}^2 A^H_{1/2}(\tau_{l'}) + \frac{v_H v_\Phi}{2 m_{H^\pm}^2} \lambda_{h_\Phi H^+ H^-} A^H_0(\tau_{H^\pm}) \right|^2, \tag{22}$$

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

$$A^H_{1/2}(\tau) = 2[\tau + (\tau - 1) f(\tau)]/\tau, \tag{23}$$

$$A^H_0(\tau) = -[\tau - f(\tau)]/\tau^2, \tag{24}$$

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

$$f(x) = \begin{cases} \arcsin^2 \sqrt{x} & , \text{ for } x \leq 1; \\ -\dfrac{1}{4}\left[\log \dfrac{1+\sqrt{1-1/x}}{1-\sqrt{1-1/x}} - i\pi\right]^2 & , \text{ for } x > 1. \end{cases} \tag{25}$$

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

$$\lambda_{h_\Phi H^+ H^-} = \tilde{\lambda}_1 \sin^2\beta + \tilde{\lambda}_2 \cos^2\beta + \sqrt{2}\mu_\Phi \sin\beta\cos\beta / v_\Phi, \tag{26}$$

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 $\mathcal{O}(10)$ 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 GeV. On the other hand, the decay width of $h_\Phi \to \gamma\gamma$ is at most $\mathcal{O}(0.1)$ GeV even for large Yukawa coupling and small m_f. For $y = 1$, the decay width is less than 5×10^{-3} GeV in the entire region of m_f. Without extra decay channels of h_Φ, we cannot achieve $O(10)$ GeV decay width for h_Φ for large m_f ($\gtrsim 500$ GeV).

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

$$\Gamma[h_\Phi \to Z\gamma] = \frac{\alpha m_W^2 m_{h_\Phi}^3}{128\pi^2 v_H^2 v_\Phi^2}\left(1 - \frac{m_Z^2}{m_{h_\Phi}^2}\right)^3 \left|\sum_{f'} N_{f'} \frac{Q_{f'} v_{f'}}{c_w} A_{1/2}^H(\tau_{f'}, \lambda_{f'})\right|^2, \tag{27}$$

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

$$A_{1/2}^H(\tau, \lambda) = I_1(\tau^{-1}, \lambda^{-1}) - I_2(\tau^{-1}, \lambda^{-1}), \tag{28}$$

where

$$I_1(x, y) = \frac{xy}{2(x-y)} + \frac{x^2 y^2}{2(x-y)^2}\left[f(x^{-1}) - f(y^{-1})\right] + \frac{x^2 y}{(x-y)^2}\left[g(x^{-1}) - g(y^{-1})\right] \tag{29}$$

$$I_2(x, y) = -\frac{xy}{2(x-y)}\left[f(x^{-1}) - f(y^{-1})\right] \tag{30}$$

with

$$g(x) = \begin{cases} \sqrt{x^{-1} - 1}\arcsin\sqrt{x} & , \text{ for } x \geq 1, \\ \dfrac{\sqrt{1 - x^{-1}}}{2}\left[\log \dfrac{1+\sqrt{1-1/x}}{1-\sqrt{1-1/x}} - i\pi\right]^2 & , \text{ for } x < 1, \end{cases} \tag{31}$$

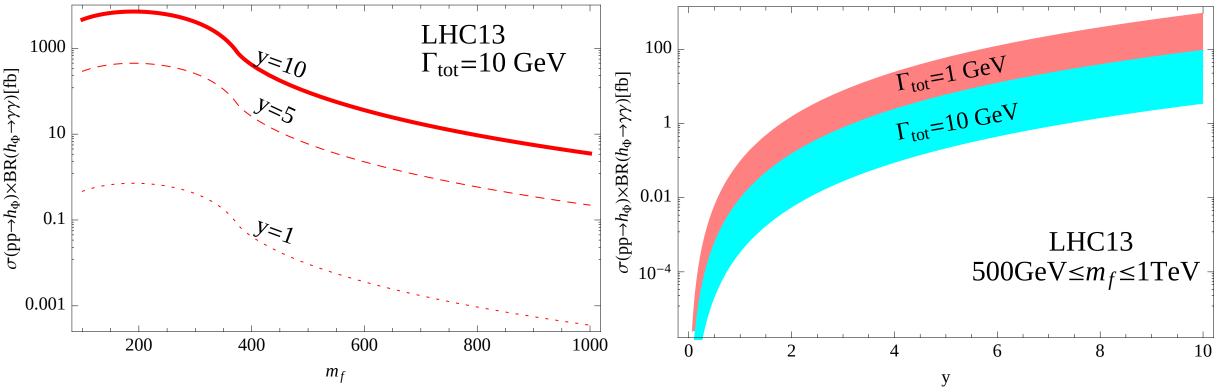

Figure 1: (Left): m_f vs. the diphoton signal via the gluon-gluon fusion at the LHC@13TeV. The total decay widths is fixed at $\Gamma_{\rm tot} = 10$ GeV. (Right): y vs. the diphoton signal for different values of the $\Gamma_{\rm tot}$ and for 500 GeV $\leq m_f \leq$ 1 TeV.

and $f(x)$ is defined in Eq. (23).

In general, h_Φ has interaction with the SM gauge bosons through the mixing between h_Φ and $h_{1,2}$ and also through the loops involving extra fermions. 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 ($\Gamma_{\rm tot}$) of h_Φ could be given by

$$\Gamma_{\rm tot} = \Delta\Gamma + \Gamma[h_\Phi \to gg] + \Gamma[h_\Phi \to \gamma\gamma] + \Gamma[h_\Phi \to \gamma Z]. \tag{32}$$

The diphoton excess requires $\mathcal{O}(10)$-GeV $\Gamma_{\rm tot}$ and $\mathcal{O}(10)$ fb diphoton signal at $\sqrt{s} = 13$ TeV. This means that large $\Delta\Gamma$ is necessary to reproduce the excess. In Fig. 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 GeV, which can be readily achieved by allowing the invisible decay of h_Φ into a pair of DM particles (see Fig. 2 and the related discussions). In the right, $\Gamma_{\rm tot} = 1$ GeV (pink) and $\Gamma_{\rm tot} = 10$ GeV (cyan) are shown, when m_f is between 500 GeV and 1 TeV. Note that we need large Yukawa coupling $y \approx 5 - 10$ for $m_f > 400$ 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 GeV, we still need large Yukawa coupling, as we see in the right panel.

Now, we discuss the detail of $\Delta\Gamma$. First of all, h_Φ can decay into $ZZ, ZZ_H, Z_H Z_H$, if there exist the mixing between $\hat{Z}$ and $\hat{Z}_H$ bosons. But it will be suppressed by small gauge coupling $g_H \lesssim O(0.1)$ and the small $Z - Z_H$ mixing. Therefore we will ignore $h_\Phi \to ZZ, ZZ_H, Z_H Z_H$.

Next, we consider the h_Φ decay into two scalar bosons if kinematically allowed. The decay width for $h_\Phi \to s_i s_j$ ($s_i = h, H, H^\pm, A$) are given by

$$\Gamma[h_\Phi \to s_i s_i] = \frac{\lambda^2_{h_\Phi s_i s_i}}{32\pi} \frac{v_\Phi^2}{m_{h_\Phi}} \sqrt{1-4x_i}, \tag{33}$$

$$\Gamma[h_\Phi \to s_i s_j] = \frac{\lambda^2_{h_\Phi s_i s_j}}{16\pi} \frac{v_\Phi^2}{m_{h_\Phi}} \lambda^{1/2}(1, x_i, x_j) (\text{for } i \neq j) \ , \tag{34}$$

$$\Gamma[h_\Phi \to \psi_X \psi_X] = \frac{y_X^2 m_{h_\Phi}}{8\pi} \{1 - 4(m^2_{\psi_X}/m^2_{h_\Phi})\}^{\frac{3}{2}}, \tag{35}$$

where $x_i = m^2_{s_i}/m^2_{h_\Phi}$ 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 for ψ_X.

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

$$\lambda_{h_\Phi hh} = \tilde{\lambda}_1 \sin^2\alpha + \tilde{\lambda}_2 \cos^2\alpha + \frac{\mu_\Phi}{\sqrt{2} v_\Phi} \sin\alpha\cos\alpha, \tag{36}$$

$$\lambda_{h_\Phi HH} = \tilde{\lambda}_1 \cos^2\alpha + \tilde{\lambda}_2 \sin^2\alpha - \frac{\mu_\Phi}{\sqrt{2} v_\Phi} \sin\alpha\cos\alpha, \tag{37}$$

$$\lambda_{h_\Phi hH} = -\tilde{\lambda}_1 \sin 2\alpha + \tilde{\lambda}_2 \cos 2\alpha - \frac{\mu_\Phi}{\sqrt{2} v_\Phi} \cos 2\alpha, \tag{38}$$

$$\begin{aligned} \lambda_{h_\Phi AA} &= \frac{v_\Phi^2}{v_\Phi^2 + (v_H \sin\beta\cos\beta)^2} \\ &\times \left(\tilde{\lambda}_1 \sin^2\beta + \tilde{\lambda}_2 \cos^2\beta + \frac{\mu_\Phi}{\sqrt{2} v_\Phi} \sin\beta\cos\beta + 2\lambda_\Phi v_H \sin\beta\cos\beta \right), \end{aligned} \tag{39}$$

where A is the pseudoscalar boson. The coupling $\lambda_{h_\Phi H^+ H^-}$ has been defined in Eq. (23) in the context of $h_\Phi \to \gamma\gamma$.

In addition, we find a dark matter candidate among the natural fermions [18]. Assuming that the Yukawa couplings are flavor-independent, we can explicitly calculate the DM mass and the Yukawa coupling with h_Φ. Fig. 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 GeV. In this case, if the total decay width of the 750 GeV excess is confirmed to be about 45 GeV, then other decay channels like $h_\Phi \to hH, HH, AA$ must be comparable to or dominant over the invisible decay † (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

†The charged scalar mass is constrained by $B \to X_s\gamma$ in type-II Higgs doublet model, and should satisfy $m_{H^\pm} \gtrsim 500$ GeV [26]. Therefore $h_\Phi \to H^+H^-$ is kinematically forbidden in our model.

there is a relation between the Z_H mass and m_f. Approximately, it can be evaluated as

$$M_{Z_H} \approx \frac{g_H}{y} m_f. \tag{40}$$

The diphoton excess suggests $\mathcal{O}(1)$ y, and then M_{Z_H} is at least $\mathcal{O}(100)$ GeV, because g_H could not be $\mathcal{O}(1)$ to evade the stringent bound from the dijet signal. Another stung constraint is from ρ parameter, as discussed in Ref. [18]. In this scenario, Z_H is light, so the coupling should be small. As discussed in Ref. [17–19], the ρ parameter is deviated from 1 at the tree-level, because of Z-Z_H mixing. The bound is roughly estimated as [19]

$$\frac{g_H}{g_Z} \frac{M_Z^2}{|M_{Z_H}^2 - M_Z^2|} \lesssim 0.004. \tag{41}$$

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 [27]. By imposing

$$\sigma(gg \to h_\Phi \to gg) \lesssim 2 \text{ pb}, \tag{42}$$

we find that the mass of the exotic quarks should be larger than 400 GeV for the Yukawa coupling $y = 5$ and 600 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 [28]:

$$\sigma(gg \to h_\Phi \to hh) \lesssim 10 \text{ fb}. \tag{43}$$

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 (43) 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 [29] and ZZ production [30] are

$$\sigma(gg \to h_\Phi \to WW) \lesssim 40 \text{ fb}, \tag{44}$$

$$\sigma(gg \to h_\Phi \to ZZ) \lesssim 10 \text{ fb}, \tag{45}$$

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) 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 [31]

$$\sigma(gg \to h_\Phi \to Z\gamma) \lesssim 3.8 \text{ fb}. \tag{46}$$

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 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 trangle 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. [5]. For more definitive conclusion on this issue, we have to perform more detailed analysis.

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,

$$\begin{aligned} q\bar{q}^{'} &\to q\bar{q}^{'} + Z_H Z_H, \quad \text{followed by} \\ Z_H Z_H &\to h_\Phi \to X X^\dagger, \end{aligned}$$

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,

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

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. (33) and (34) are determined by the dimensionless couplings in the Higgs potential, Eqs. (36)–(39). In principle, these extra decay could be sizable up to $\Delta\Gamma \approx \mathcal{O}(10)$ 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.

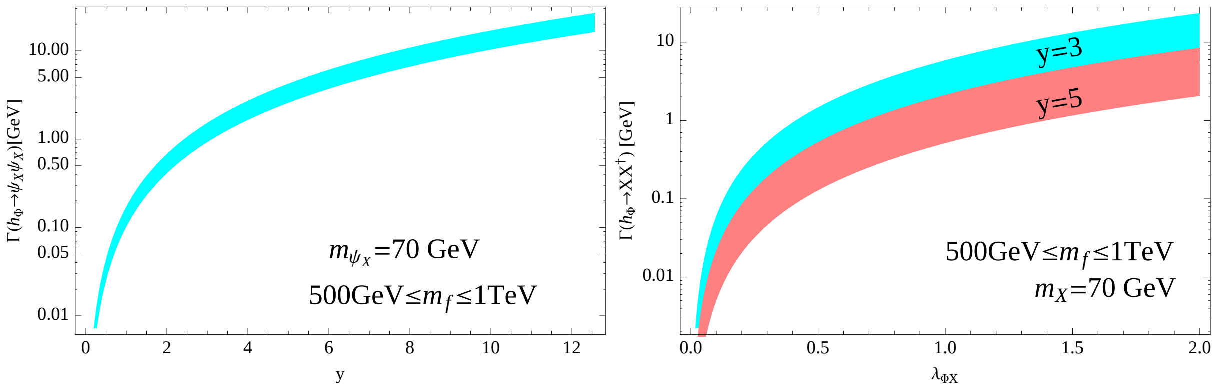

Figure 2: y vs. invisible decay width of h_Φ (GeV) in the fermionic DM scenario (left) and scalar DM scenario (right).The vector-like fermion mass is between 500 GeV and 1 TeV on the cyan and pink bands. The dark matter masses are 70 GeV in the both cases.

4 Summary

In this paper we interpret the recently reported diphoton excess at 750 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 $\mathrm{U(1)}'$ so that their masses arise entirely from spontaneous breaking of $\mathrm{U(1)}'$ 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 $\mathrm{U(1)}'$ 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 $O(10)$ GeV decay width, making one of the dominant decay channels of h_Φ.

If the diphoton excess at 750 GeV with large decay width ~ 45 GeV is confirmed in the next LHC run, our model predicts that there should be new vectorlike quarks and leptons around $\sim O(1)$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. 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 GeV diphoton excess will be presented elsewhere.

Acknowledgments

This work is supported in part by National Research Foundation of Korea (NRF) Research Grant NRF-2015R1A2A1A05001869, and by the NRF grant funded by the Korea government (MSIP) (No. 2009-0083526) through Korea Neutrino Research Center at Seoul National University (P.K.). The work of Y.O. is supported by Grant-in-Aid for Scientific research from the Ministry of Education, Science, Sports, and Culture (MEXT), Japan, No. 23104011. The work of C.Y. was supported by the Ministry of Science and Technology (MoST) of Taiwan under grant number 101-2112-M-001-005-MY3.

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