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Nuclear Physics B 909 (2016) 122–134

NUCLEAR PHYSICS B

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Gauged Two Higgs Doublet Model confronts the LHC 750 GeV diphoton anomaly

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Received 19 January 2016; received in revised form 5 April 2016; accepted 1 May 2016

Available online 4 May 2016

Editor: Hong-Jian He

Abstract

In light of the recent 750 GeV diphoton anomaly observed at the LHC, we study the possibility of accommodating the deviation from the standard model prediction based on the recently proposed Gauged Two Higgs Doublet Model. The model embeds two Higgs doublets into a doublet of a non-abelian gauge group $SU(2)_H$, while the standard model $SU(2)_L$ right-handed fermion singlets are paired up with new heavy fermions to form $SU(2)_H$ doublets, and $SU(2)_L$ left-handed fermion doublets are singlets under $SU(2)_H$. An $SU(2)_H$ scalar doublet, which provides masses to the new heavy fermions as well as the $SU(2)_H$ gauge bosons, can be produced via gluon fusion and subsequently decays into two photons with the new fermions circulating the triangle loops to account for the deviation from the standard model prediction.

1. Introduction

Recent results from LHC [1–3] exhibit an intriguing anomaly on the diphoton channel at the scale around 750 GeV. Numerous attempts [4–81] have been put forward to explain the excess, while Refs. [14,43,57] are based on two Higgs doublet models, similar to this work.

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http://dx.doi.org/10.1016/j.nuclphysb.2016.05.002
0550-3213/

In Ref. [78], a combined result from run I and II gives a cross section $\sigma(pp \to X \to \gamma\gamma) \sim \mathcal{O}(6)$ fb for a scalar particle X with mass around 750 GeV. In this paper, we will show that the newly proposed Gauged Two Higgs Doublet Model [82] (G2HDM) is able to provide a cross section with such magnitude.

G2HDM contains additional $SU(2)_H \times U(1)_X$ gauge symmetry, in which H_1 (identified as the Standard Model (SM) Higgs doublet) and H_2 comprise an $SU(2)_H$ doublet such that the two-doublet potential is as simple as the SM Higgs potential with just a quadratic mass term plus a quartic term. The cost to pay is to include additional scalars: one $SU(2)_H$ triplet Δ_H and one $SU(2)_H$ doublet Φ_H (that are all singlets under the SM gauge groups) with their vacuum expectation values (vevs) supplying masses to the $SU(2)_H \times U(1)_X$ gauge bosons. Moreover, the vev of the triplet induces the SM Higgs vev, breaking $SU(2)_L \times U(1)_Y$ down to $U(1)_Q$, while H_2 does not obtain any vev and the neutral component of H_2 could be a dark matter (DM) candidate, whose stability is protected by the $SU(2)_H$ gauge symmetry and Lorentz invariance, without resorting to an ad-hoc Z_2 symmetry. In order to write down $SU(2)_H \times U(1)_X$ invariant Yukawa couplings, we introduce heavy $SU(2)_L$ singlet Dirac fermions, the right-handed component of which is paired up with the SM right-handed fermions to comprise $SU(2)_H$ doublets. In this setup, the model is anomaly-free regarding all gauge groups involved.

In this work, we focus on the role of ϕ_2 which is a physical component in Φ_H and whose vev $\langle\phi_2\rangle = v_\Phi$ gives masses to the new heavy fermions. Since it couples to new colored fermions, it can be produced radiatively via gluon fusion and also decay radiatively into a pair of photons with the heavy charged fermions in loops. We will demonstrate that ϕ_2 can be a good candidate if LHC eventually could confirm the diphoton anomaly. Moreover, the observed width of the bump can be simply obtained from ϕ_2 decay into the additional fermions with $\mathcal{O}(1)$ Yukawa couplings.

The paper is organized as follows. First, we briefly discuss the G2HDM in Section 2 restraining ourselves only to those aspects most relevant to $\gamma\gamma$ mode. Next, in Section 3 we compute the diphoton cross section through ϕ_2 exchange and the partial decay width of ϕ_2 into the new heavy fermions. In Section 4, we briefly comment on implications of such the new heavy fermions in terms of collider searches, electron and muon magnetic dipole moment measurements, and the electroweak precision test data. Finally, we conclude in Section 5.

2. G2HDM setup

In this Section, we review the G2HDM (cf. Ref. [82]) with the particle content summarized in Table 1.

For the scalar sector, we have two Higgs doublets, H_1 and H_2, where H_1 is identified as the SM Higgs doublet and H_2 (with the same hypercharge $Y = 1/2$ as H_1) is the extra $SU(2)_L$ doublet. H_1 and H_2 transform as a doublet $H = (H_1\ H_2)^T$ under the additional gauge group $SU(2)_H \times U(1)_X$ with $U(1)_X$ charge $X(H) = 1$. Besides the doublet H, we also introduce $SU(2)_H$ triplet and doublet, Δ_H and Φ_H, which are *singlets* under $SU(2)_L$. The Higgs potential invariant under both $SU(2)_L \times U(1)_Y$ and $SU(2)_H \times U(1)_X$ can be written down easily as

$$V(H, \Delta_H, \Phi_H) = V(H) + V(\Phi_H) + V(\Delta_H) + V_{\rm mix}(H, \Delta_H, \Phi_H) \ , \tag{1}$$

with

$$\begin{aligned} V(H) &= \mu_H^2 H^\dagger H + \lambda_H \left(H^\dagger H\right)^2 \ , \\ &= \mu_H^2 \left(H_1^\dagger H_1 + H_2^\dagger H_2\right) + \lambda_H \left(H_1^\dagger H_1 + H_2^\dagger H_2\right)^2 \ , \end{aligned} \tag{2}$$

Table 1
Matter field contents and their quantum number assignments in G2HDM.

Matter fields	$SU(3)_C$	$SU(2)_L$	$SU(2)_H$	$U(1)_Y$	$U(1)_X$
$Q_L = (u_L \;\; d_L)^T$	3	2	1	1/6	0
$U_R = \left(u_R \;\; u_R^H\right)^T$	3	1	2	2/3	1
$D_R = \left(d_R^H \;\; d_R\right)^T$	3	1	2	−1/3	−1
$L_L = (\nu_L \;\; e_L)^T$	1	2	1	−1/2	0
$N_R = \left(\nu_R \;\; \nu_R^H\right)^T$	1	1	2	0	1
$E_R = \left(e_R^H \;\; e_R\right)^T$	1	1	2	−1	−1
χ_u	3	1	1	2/3	0
χ_d	3	1	1	−1/3	0
χ_ν	1	1	1	0	0
χ_e	1	1	1	−1	0
$H = (H_1 \;\; H_2)^T$	1	2	2	1/2	1
$\Delta_H = \begin{pmatrix} \Delta_3/2 & \Delta_p/\sqrt{2} \\ \Delta_m/\sqrt{2} & -\Delta_3/2 \end{pmatrix}$	1	1	3	0	0
$\Phi_H = (\Phi_1 \;\; \Phi_2)^T$	1	1	2	0	1

$$
\begin{aligned}
V(\Phi_H) &= \mu_\Phi^2 \Phi_H^\dagger \Phi_H + \lambda_\Phi \left(\Phi_H^\dagger \Phi_H\right)^2 \; , \\
&= \mu_\Phi^2 \left(\Phi_1^* \Phi_1 + \Phi_2^* \Phi_2\right) + \lambda_\Phi \left(\Phi_1^* \Phi_1 + \Phi_2^* \Phi_2\right)^2 \; , \qquad (3) \\
V(\Delta_H) &= -\mu_\Delta^2 \mathrm{Tr}\left(\Delta_H^\dagger \Delta_H\right) + \lambda_\Delta \left(\mathrm{Tr}\left(\Delta_H^\dagger \Delta_H\right)\right)^2 \; , \\
&= -\mu_\Delta^2 \left(\frac{1}{2}\Delta_3^2 + \Delta_p \Delta_m\right) + \lambda_\Delta \left(\frac{1}{2}\Delta_3^2 + \Delta_p \Delta_m\right)^2 \; , \qquad (4)
\end{aligned}
$$

and finally the mixed term

$$
\begin{aligned}
V_{\rm mix}\left(H, \Delta_H, \Phi_H\right) = & + M_{H\Delta}\left(H^\dagger \Delta_H H\right) - M_{\Phi\Delta}\left(\Phi_H^\dagger \Delta_H \Phi_H\right) \\
& + \lambda_{H\Delta}\left(H^\dagger H\right) \mathrm{Tr}\left(\Delta_H^\dagger \Delta_H\right) + \lambda_{H\Phi}\left(H^\dagger H\right)\left(\Phi_H^\dagger \Phi_H\right) \\
& + \lambda_{\Phi\Delta}\left(\Phi_H^\dagger \Phi_H\right) \mathrm{Tr}\left(\Delta_H^\dagger \Delta_H\right) \; , \\
= & + M_{H\Delta}\left(\frac{1}{\sqrt{2}} H_1^\dagger H_2 \Delta_p + \frac{1}{2} H_1^\dagger H_1 \Delta_3 + \frac{1}{\sqrt{2}} H_2^\dagger H_1 \Delta_m - \frac{1}{2} H_2^\dagger H_2 \Delta_3\right) \\
& - M_{\Phi\Delta}\left(\frac{1}{\sqrt{2}} \Phi_1^* \Phi_2 \Delta_p + \frac{1}{2} \Phi_1^* \Phi_1 \Delta_3 + \frac{1}{\sqrt{2}} \Phi_2^* \Phi_1 \Delta_m - \frac{1}{2} \Phi_2^* \Phi_2 \Delta_3\right) \\
& + \lambda_{H\Delta}\left(H_1^\dagger H_1 + H_2^\dagger H_2\right)\left(\frac{1}{2}\Delta_3^2 + \Delta_p \Delta_m\right) \\
& + \lambda_{H\Phi}\left(H_1^\dagger H_1 + H_2^\dagger H_2\right)\left(\Phi_1^* \Phi_1 + \Phi_2^* \Phi_2\right) \\
& + \lambda_{\Phi\Delta}\left(\Phi_1^* \Phi_1 + \Phi_2^* \Phi_2\right)\left(\frac{1}{2}\Delta_3^2 + \Delta_p \Delta_m\right) \; , \qquad (5)
\end{aligned}
$$

where

$$\Delta_H = \begin{pmatrix} \Delta_3/2 & \Delta_p/\sqrt{2} \\ \Delta_m/\sqrt{2} & -\Delta_3/2 \end{pmatrix} \text{ with } \Delta_m = \left(\Delta_p\right)^* \text{ and } (\Delta_3)^* = \Delta_3 \,, \tag{6}$$

and $\Phi_H = (\Phi_1 \;\; \Phi_2)^T$.

Note that the quadratic terms of H_1 and H_2 have the following coefficients

$$\mu_H^2 \mp \frac{1}{2} M_{H\Delta} \cdot v_\Delta + \frac{1}{2}\lambda_{H\Delta} \cdot v_\Delta^2 + \frac{1}{2}\lambda_{H\Phi} \cdot v_\Phi^2 \,, \tag{7}$$

respectively. As a result even with a positive μ_H^2, H_1 can still develop a vev $(0\, v/\sqrt{2})^T$ breaking $SU(2)_L$ provided that the second term is dominant, while H_2 remains zero vev. In other words, electroweak symmetry breaking is triggered by the $SU(2)_H$ breaking.

To facilitate electroweak symmetry breaking spontaneously, it is convenience to parametrize the scalars as

$$H_1 = \begin{pmatrix} G^+ \\ \frac{v+h}{\sqrt{2}} + iG^0 \end{pmatrix} \,, \quad \Phi_H = \begin{pmatrix} G_H^p \\ \frac{v_\Phi + \phi_2}{\sqrt{2}} + iG_H^0 \end{pmatrix} \,, \quad \Delta_H = \begin{pmatrix} \frac{-v_\Delta + \delta_3}{2} & \frac{1}{\sqrt{2}}\Delta_p \\ \frac{1}{\sqrt{2}}\Delta_m & \frac{v_\Delta - \delta_3}{2} \end{pmatrix} \tag{8}$$

and $H_2 = (H_2^+ \; H_2^0)^T$. Here v, v_Φ and v_Δ are vevs to be decided by minimizing the potential. $\Psi_G \equiv \{G^+, G^3, G_H^p, G_H^0\}$ are Goldstone bosons, to be eaten by the longitudinal components of W^+, W^3, W^p, W'^3 respectively, while $\Psi \equiv \{h, H_2, \Phi_1, \phi_2, \delta_3, \Delta_p\}$ are the physical fields.

Nonzero vevs v, v_Φ and v_Δ will induce the mixing among the scalars, leading to two mass matrices. In this work, the relevant mass matrix in the basis of $\{h, \delta_3, \phi_2\}$ is given by

$$\mathcal{M}_0^2 = \begin{pmatrix} 2\lambda_H v^2 & \frac{v}{2}\left(M_{H\Delta} - 2\lambda_{H\Delta} v_\Delta\right) & \lambda_{H\Phi} v v_\Phi \\ \frac{v}{2}\left(M_{H\Delta} - 2\lambda_{H\Delta} v_\Delta\right) & \frac{1}{4 v_\Delta}\left(8\lambda_\Delta v_\Delta^3 + M_{H\Delta} v^2 + M_{\Phi\Delta} v_\Phi^2\right) & \frac{v_\Phi}{2}\left(M_{\Phi\Delta} - 2\lambda_{\Phi\Delta} v_\Delta\right) \\ \lambda_{H\Phi} v v_\Phi & \frac{v_\Phi}{2}\left(M_{\Phi\Delta} - 2\lambda_{\Phi\Delta} v_\Delta\right) & 2\lambda_\Phi v_\Phi^2 \end{pmatrix} . \tag{9}$$

To simplify the diphoton excess analysis below, we focus on the simplest but representative scenario where all off-diagonal terms vanish by choosing

$$\lambda_{H\Phi} = 0 \,, \; M_{H\Delta} = 2\lambda_{H\Delta} v_\Delta \,, \; M_{\Phi\Delta} = 2\lambda_{\Phi\Delta} v_\Delta \,, \tag{10}$$

and the scalar masses become

$$m_h^2 = 2\lambda_H v^2 \,, \; m_{\delta_3}^2 = \frac{1}{2}\left(4\lambda_\Delta v_\Delta^2 + \lambda_{H\Delta} v^2 + \lambda_{\Phi\Delta} v_\Phi^2\right) \,, \; m_{\phi_2}^2 = 2\lambda_\Phi v_\Phi^2 \,, \tag{11}$$

where the value of λ_H is exactly the same as in the SM. In this scenario, there is no mixing among h, δ_3, ϕ_2 [1] and the scalar ϕ_2 is responsible for the diphoton excess as we shall see below.

Next, the fermion sector together with new Yukawa couplings will be discussed. By virtue of the additional gauge group $SU(2)_H$, new heavy fermions have to be included but there are various ways to implement the idea. We, however, stick to the simplest realization: the heavy fermions together with the SM right-handed fermions form $SU(2)_H$ doublets, while the SM left-handed doublets are singlets under $SU(2)_H$. We begin with the quark sector. In the simplest realization, one can make the quark $SU(2)_L$ doublet, Q_L, an $SU(2)_H$ singlet and incorporate

[1] Therefore, subtleties from the scalar mixing, for example, the impact on electroweak vacuum stability [83] will not be discussed here.

extra $SU(2)_L$ singlets u_R^H and d_R^H which together with the SM right-handed quarks u_R and d_R, respectively, form $SU(2)_H$ doublets: $U_R^T = (u_R\ u_R^H)_{2/3}$ and $D_R^T = (d_R^H\ d_R)_{-1/3}$, where the subscript denotes hypercharge. As a consequence, we have Yukawa couplings

$$\begin{aligned}\mathcal{L}_{\text{Yuk}} &\supset y_d \bar{Q}_L \left(D_R \cdot H\right) + y_u \bar{Q}_L \left(U_R \cdot \tilde{\tilde{H}}\right) + \text{H.c.},\\ &= y_d \bar{Q}_L \left(d_R^H H_2 - d_R H_1\right) - y_u \bar{Q}_L \left(u_R \tilde{H}_1 + u_R^H \tilde{H}_2\right) + \text{H.c.},\end{aligned} \tag{12}$$

where "·" refers to $SU(2)_H$ multiplication[2] and $\tilde{\tilde{H}} \equiv (\tilde{H}_2\ -\tilde{H}_1)^T$ with $\tilde{H}_{1,2} = i\tau_2 H_{1,2}^*$ transforms as 2 under $SU(2)_H$. After the electroweak symmetry breaking $\langle H_1 \rangle \neq 0$, u and d obtain their masses but u^H and d^H remain massless since H_2 does not get a vev.

To provide masses to the additional species, we make use of the $SU(2)_H$ scalar doublet $\Phi_H = (\Phi_1\ \Phi_2)^T$, which is neutral under $SU(2)_L$, and left-handed $SU(2)_{L,H}$ singlets χ_u and χ_d as

$$\begin{aligned}\mathcal{L}_{\text{Yuk}} &\supset -y'_d \overline{\chi}_d \left(D_R \cdot \Phi_H\right) + y'_u \overline{\chi}_u \left(U_R \cdot \tilde{\Phi}_H\right) + \text{H.c.},\\ &= -y'_d \overline{\chi}_d \left(d_R^H \Phi_2 - d_R \Phi_1\right) - y'_u \overline{\chi}_u \left(u_R \Phi_1^* + u_R^H \Phi_2^*\right) + \text{H.c.},\end{aligned} \tag{13}$$

in which Φ_H has $Y = 0$, $Y(\chi_u) = Y(U_R) = 2/3$ and $Y(\chi_d) = Y(D_R) = -1/3$ with $\tilde{\Phi}_H = (\Phi_2^*\ -\Phi_1^*)^T$. With $\langle \Phi_2 \rangle = v_\Phi/\sqrt{2}$, u^H (χ_u) and d^H (χ_d) obtain masses $y'_u v_\Phi/\sqrt{2}$ and $y'_d v_\Phi/\sqrt{2}$, respectively. Notice that both v_Δ and v_Φ contribute to the $SU(2)_H$ gauge boson masses.

The lepton sector mimics the quark sector as

$$\begin{aligned}\mathcal{L}_{\text{Yuk}} &\supset y_e \bar{L}_L \left(E_R \cdot H\right) + y_\nu \bar{L}_L \left(N_R \cdot \tilde{H}\right) - y'_e \overline{\chi}_e \left(E_R \cdot \Phi_H\right) + y'_\nu \overline{\chi}_\nu \left(N_R \cdot \tilde{\Phi}_H\right) + \text{H.c.},\\ &= y_e \bar{L}_L \left(e_R^H H_2 - e_R H_1\right) - y_\nu \bar{L}_L \left(\nu_R \tilde{H}_1 + \nu_R^H \tilde{H}_2\right)\\ &\quad - y'_e \overline{\chi}_e \left(e_R^H \Phi_2 - e_R \Phi_1\right) - y'_\nu \overline{\chi}_\nu \left(\nu_R \Phi_1^* + \nu_R^H \Phi_2^*\right) + \text{H.c.},\end{aligned} \tag{14}$$

in which $E_R^T = (e_R^H\ e_R)_{-1}$ and $N_R^T = (\nu_R\ \nu_R^H)_0$ where ν_R and ν_R^H correspond to the right-handed neutrino and the $SU(2)_H$ partner of it respectively, while χ_e and χ_ν are $SU(2)_{L,H}$ singlets with $Y(\chi_e) = -1$ and $Y(\chi_\nu) = 0$. Similarly all SM leptons and their heavy counterparts will obtain masses from $\langle H_1 \rangle$ and $\langle \Phi_2 \rangle$.

As mentioned above, because ϕ_2 (a member of Φ_H) couples to the new heavy fermions, it can be radiatively produced via loops of the new colored particles and radiatively decays into the diphoton final state via loops of the new charged particles to accommodate the observed bump. On the other hand, although ϕ_2 is a singlet under the SM gauge group, it does couple to SM fermions and gauge bosons at tree level via the h–ϕ_2 mixing. That is the reason why we work in the zero mixing limit to evade direct search bounds from, for instance, dijet or dilepton channels. Note that there are no excesses in the ZZ, dijet or dilepton channels near the invariant mass of 750 GeV.

3. Diphoton anomaly

Equipped with the basics of G2HDM, we are now in a position to calculate the diphoton cross section via ϕ_2 exchange. The cross section at the ϕ_2-resonance can be well approximated by [84]

[2] For 2-dimensional $SU(2)_H$ spinors A and B, $A \cdot B = \epsilon_{ij} A^i B^j$.

$$\sigma\left(gg \to \phi_2 \to \gamma\gamma\right) = \frac{\pi^2}{8s\, m_{\phi_2}\, \Gamma_{\phi_2}} f_{gg}\left(\frac{m_{\phi_2}}{\sqrt{s}}\right) \Gamma\left(\phi_2 \to gg\right) \Gamma\left(\phi_2 \to \gamma\gamma\right), \tag{15}$$

with the center of mass energy $\sqrt{s} = 13$ TeV and the integral of the parton (gluon in this case) distribution function product

$$f_{gg} = \int_{m_{\phi_2}^2/s}^{1} \frac{dx}{x} g\left(x, \mu^2\right) g\left(\frac{m_{\phi_2}^2}{sx}, \mu^2\right) = 2141.7, \tag{16}$$

evaluated at the scale $\mu = m_{\phi_2}$, using MSTW2008NNLO [85] and the value is consistent with Ref. [15]. The partial decay width of ϕ_2 into a heavy fermion and antifermion in the presence of a Yukawa term, $y'_f \phi_2 \bar{f} f/\sqrt{2}$, that also gives a mass m_f to the heavy fermion because of $\langle\phi_2\rangle = v_\Phi$, reads

$$\Gamma\left(\phi_2 \to f\bar{f}\right) = N_c \frac{y_f'^2\, m_{\phi_2}}{16\pi}\left(1 - 4\frac{m_f^2}{m_{\phi_2}^2}\right)^{3/2}, \tag{17}$$

where $N_c = 3$ for heavy colored particles while $N_c = 1$ for heavy leptons.

The partial decay width of ϕ_2 into diphoton mediated by heavy fermions is [86–88]

$$\Gamma\left(\phi_2 \to \gamma\gamma\right) = \frac{\alpha^2\, m_{\phi_2}^3}{256\, v_\Phi^2\, \pi^3}\left|\sum_f N_c Q_f^2 A_{1/2}^H(\tau_f)\right|^2, \tag{18}$$

where $\tau_f = m_{\phi_2}^2/4m_f^2$ with

$$A_{1/2}^H(\tau) = 2[\tau + (\tau - 1) f(\tau)]\, \tau^{-2}, \tag{19}$$

and the function $f(\tau)$ is defined as

$$f(\tau) = \begin{cases} \arcsin^2\sqrt{\tau} & , \text{ for } \tau \le 1\,; \\ -\dfrac{1}{4}\left[\log\dfrac{1+\sqrt{1-\tau^{-1}}}{1-\sqrt{1-\tau^{-1}}} - i\pi\right]^2 & , \text{ for } \tau > 1\,. \end{cases} \tag{20}$$

On the other hand, the partial decay width of ϕ_2 into 2 gluons mediated by colored heavy fermions is [86–88]

$$\Gamma\left(\phi_2 \to gg\right) = \frac{\alpha_s^2\, m_{\phi_2}^3}{72\, v_\Phi^2\, \pi^3}\left|\sum_f \frac{3}{4} A_{1/2}^H(\tau_f)\right|^2. \tag{21}$$

In our model, there are 6 heavy colored Dirac fermions, including 3 generations of up-type and down-type heavy quarks (with electric charge of 2/3 and 1/3, respectively) which contribute in $\Gamma\left(\phi_2 \to gg\right)$ while for $\Gamma\left(\phi_2 \to \gamma\gamma\right)$ there are additional 3 heavy charged leptons with one unit of electric charge in addition to the heavy quarks. From the CMS run I and CMS+ATLAS run II diphoton data combined, the best fit value for the diphoton cross section is 6.2 ± 1.0 femtobarn [78]. It implies in units of GeV^{-2}

$$\sigma\left(gg \to \phi_2 \to \gamma\gamma\right) = \frac{f_{gg}\left(\frac{m_{\phi_2}}{\sqrt{s}}\right)\pi^2}{8s} \frac{m_{\phi_2}}{\Gamma_{\phi_2}} \frac{\Gamma\left(\phi_2 \to gg\right)}{m_{\phi_2}} \frac{\Gamma\left(\phi_2 \to \gamma\gamma\right)}{m_{\phi_2}} \simeq 1.60 \times 10^{-11}, \tag{22}$$

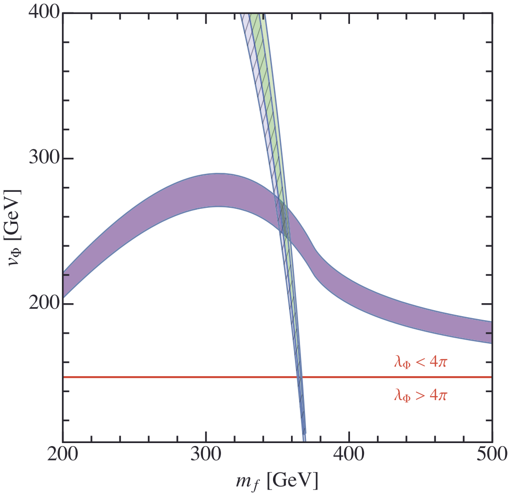

Fig. 1. The purple area on the m_f–v_Φ plane is the 1σ region which reproduces the $\gamma\gamma$ bump at the LHC. The green shaded region denotes $0.05 < \Gamma_{\phi_2}/m_{\phi_2} < 0.07$, including all neutral and charged heavy fermions, while the blue shaded region takes into account the heavy charged particles only. The red solid line marks the perturbativity limit because $m^2_{\phi_2} = 2\lambda_\Phi v^2_\Phi$. In order to reproduce the diphoton bump with the proper width, one will need fermion masses to be around 360 GeV and the vev v_Φ at 250 GeV, implying $\mathcal{O}(1)$ Yukawa couplings. (For interpretation of the references to color in this figure, the reader is referred to the web version of this article.)

i.e.,

$$1.65 \times 10^{-8} \simeq \frac{\Gamma(\phi_2 \to gg)}{m_{\phi_2}} \frac{\Gamma(\phi_2 \to \gamma\gamma)}{m_{\phi_2}}, \tag{23}$$

with $\sqrt{s} = 13$ TeV and $\Gamma_{\phi_2}/m_{\phi_2} \simeq 0.06$ [1].

In the Fig. 1, we color the 1σ region in purple on the m_f–v_Φ plane to accommodate the $\gamma\gamma$ anomaly where all heavy fermions involved are assumed to have the same mass m_f for simplicity. The green shaded region corresponds to the total decay width of ϕ_2, obtained from Eq. (17) by including all neutral and charged heavy fermions (u^H, d^H, e^H, ν^H), at the range of $0.05 < \Gamma_{\phi_2}/m_{\phi_2} < 0.07$ that is consistent with the observed resonance width [1]. By contrast, the blue shaded region denotes the total decay width of ϕ_2 with $0.05 < \Gamma_{\phi_2}/m_{\phi_2} < 0.07$, including heavy charged particles only (u^H, d^H, e^H). The red solid line corresponds to the perturbativity limit since $m^2_{\phi_2} = 2\lambda_\Phi v^2_\Phi$ in the limit of zero mixing among h, ϕ_2 and δ_3. In order to have the diphoton excess, one can see that the new fermion masses have to be around 360 GeV with the vev v_Φ at 250 GeV. However, we can also relax our assumption to allow for non-degenerate heavy fermion masses. In this case, one can still achieve the diphoton excess and the desired total decay width of ϕ_2, while the heavy charged fermion masses are not longer constrained to be around 360 GeV.

We conclude this Section by commenting on impacts of having v_Φ around 250 GeV. As discussed in Ref. [82], v_Φ is restrained to be of order TeV to avoid various constraints. Small v_Φ will induce a large mixing between the SM Z and $SU(2)_H$ Z', which can be avoided if the $SU(2)_H$ gauge coupling g_H is small. To be more clear, the mixing angle, in the limit of $g_H \ll g$, reads

$$\sin\theta_{ZZ'} \simeq -\frac{g_H}{\sqrt{g^2 + g'}}, \tag{24}$$

where g and g' are the SM $SU(2)_L$ and $U(1)_Y$ gauge coupling constants, respectively. One can in principle make g_H small to have a very small mixing, resulting in very light $SU(2)_H$ gauge bosons. On the other hand, the DM matter candidate in this case could be the new neutral lepton (ν_R^H or χ_ν), the $SU(2)_H$ W' or the neutral Higgs H_2^0, depending on the parameter space. The DM stability is protected by the $SU(2)_H$ gauge symmetry and the Lorentz invariance as demonstrated in Ref. [82].

4. Implications of a few hundred GeV heavy fermions

In this Section, we briefly comment on some of consequences of $SU(2)_H$ heavy fermions with masses of order 360 GeV, required to realize the diphoton excess. A detailed study is, however, beyond the scope of this paper and deserves a separate work.

4.1. Muon and electron magnetic dipole moment $g-2$

At one-loop level, the charged leptons (electron and muon) anomalous magnetic moment ($g_\ell - 2$) receive three additional radiative contributions[3] involving loops of W' with ℓ^H, H_2 with ℓ^H and Z' with ℓ^H, out of which the H_2 contribution can be neglected because it is highly suppressed by the corresponding small SM electron and muon Yukawa couplings and H_2 are assumed to be heavy. Taking into account the fact W' and Z' only couple to the right-handed SM fermions, the gauge boson contributions to the anomaly $a_\ell \equiv (g_\ell - 2)/2$ are [89,90]

$$
\begin{aligned}
a_l^{W'} &= \frac{g_H^2}{32\pi^2} \int_0^1 dx \, \frac{(1-x)}{r_{W'}^2 \left(r_H^2 (1-x) + \left(r_{W'}^2 - (1-x)\right) x\right)} \\
&\quad \times \left(r_H (1-x)^3 + 4 r_H r_{W'}^2 x + (1-x)^2 x - \left(r_H^2 (1-x)^2 + 2 r_{W'}^2 x (1+x) \right) \right) \\
&\simeq \frac{g_H^2}{48\pi^2} \begin{cases} \frac{12 r_H r_{W'}^2 - 9 r_{W'}^2 - 2 r_H^2}{4 r_H^2 r_{W'}^2} & \text{for } m_{\ell^H} \gg m_{W'} > m_\ell \ , \\ \frac{3 r_H - 2}{r_{W'}^2} & \text{for } m_{W'} \gg m_{\ell^H} > m_\ell \ , \end{cases}
\end{aligned}
\tag{25}
$$

and

$$
\begin{aligned}
a_l^{Z'} &= \frac{g_H^2}{32\pi^2} \int_0^1 dx \, \frac{x (1-x)^2}{(1-x)^2 + r_{Z'}^2 x} \\
&\simeq \frac{g_H^2}{64\pi^2} \begin{cases} 1 & \text{for } m_\ell \gg m_{Z'} \ , \\ \frac{2}{3 r_{Z'}^2} & \text{for } m_{Z'} \gg m_\ell \ , \end{cases}
\end{aligned}
\tag{26}
$$

where $r_H \equiv m_{\ell^H}/m_\ell$ and $r_{(W',Z')} \equiv m_{(W',Z')}/m_\ell$.

In addition, the Z–Z' mixing with the angle given in Eq. (24) also induces an extra contribution to a_l, obtained by multiplying Eq. (26) by $(\sin\theta_{ZZ'})^2$ and replacing g_H by $g/(\cos\theta_w)$, where θ_w is the Weinberg angle. In contrast, due to the quantum number assignment, W' is electrically neutral and will not mix with the SM W boson, unlike Z'. Thus, Eq. (25) is the total contribution from W'. Moreover, the W' and Z' boson masses are

[3] To simplify the analysis, we treat $U(1)_X$ as a global symmetry by setting $g_X = 0$.

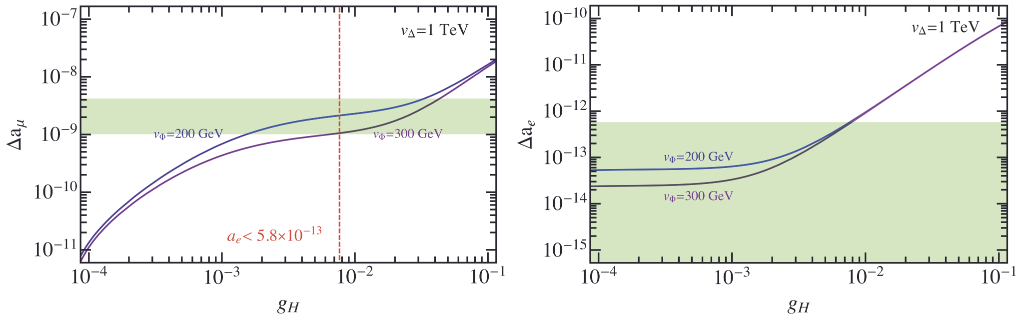

Fig. 2. Muon and electron $\Delta a_l = (g-2)_l/2$. (For interpretation of the references to color in this figure, the reader is referred to the web version of this article.)

$$m_{W'}^2 = \frac{1}{4} g_H^2 \left(v^2 + v_\Phi^2 + 4v_\Delta^2 \right) ,$$
$$m_{Z'}^2 \simeq \frac{1}{4} g_H^2 v_\Phi^2 , \quad (\text{ in the limit of } g_H \ll g, g') . \tag{27}$$

We present our results in Fig. 2 where all of Z', W' and Z–Z' mixing contributions are included. In the left-panel, with v_Δ set to 1 TeV and $m_{\ell H}$ to be 360 GeV, the green band on the g_H–Δa_μ plane corresponds to the 2σ region of the difference between the experimental value and the SM prediction [91–93], $10.1 \times 10^{-10} < a_\mu^{\rm exp} - a_\mu^{\rm SM} < 42.1 \times 10^{-10}$, the blue (purple) line refers to $v_\Phi = 200$ (300) GeV. To explain the muon anomaly Δa_μ, small values of v_Φ are preferred. The red dashed line is the limit extracted from the electron anomaly Δa_e as shown in the right panel, where the green band represents $-2.7 \times 10^{-12} < a_e^{\rm exp} - a_e^{\rm SM} < 5.8 \times 10^{-13}$ [94–97]. For $g_H \lesssim 10^{-3}$, the electron anomaly Δa_e scales as $g_H^2 m_e^2 / m_{(W',Z')}^2$, which is simply $m_e^2/v_{(\Delta,\Phi)}^2$ since $m_{(W',Z')}^2 \sim g_H^2 v_{(\Delta,\Phi)}^2$. This implies independence of Δa_e on g_H. However, for $g_H \gtrsim 10^{-2}$ it is proportional to g_H^2, since for $m_{\ell H} \sim m_{W'} \gg m_\ell$, $a_e^{W'} \sim g_H \frac{m_\ell}{m_{\ell H}}$ from Eq. (25).

4.2. Collider searches

In previous subsection, we showed that in order to accommodate the diphoton excess without contradicting the electron and muon $g-2$ measurement, the $SU(2)_H$ gauge coupling g_H is confined to be less than 10^{-2}. Thus, at the LHC the heavy fermions will be mainly produced via the 750 GeV ϕ_2 decay due to large Yukawa couplings of $\mathcal{O}(1)$ instead of being generated through W'- and Z'-exchange processes. By virtue of the $SU(2)_H$ gauge symmetry, the decay of these heavy fermions must be accompanied by the DM particle in the final state as well.

For illustration, we use τ^H as an example. It has three different decay channels, corresponding to three possible DM candidates ν^H, H_2^0 and W' in G2HDM, respectively:

$$\tau^H \to W'^p \, \tau_R \to \nu^H \, \overline{\nu_R} \, \tau_R ,$$
$$\tau^H \to H_2^0 \, \tau_L ,$$
$$\tau^H \to W'^p \, \tau_R , \tag{28}$$

where in the first channel one could have multiple leptons or jets in addition to missing energy depending on whether $\overline{\nu_R}$ decays into $\overline{\nu_L}$ and H_1 within the detector or not, while the last two channels feature one lepton plus missing transverse energy.

The energy of SM fermion τ in the final state depends on the mass difference between τ^H and the DM. If the mass splitting is too small, this may lead to very soft τ which fails to pass the event selection. The process $gg \to \phi_2 \to \overline{\tau^H}\tau^H \to$ null (DM + soft τs), which will be largely excluded by the DM mono-jet searches as pointed out in Ref. [98]. On the other hand, if the mass splitting is large enough, the final state τ is visible and the situation will require delicate study, see Ref. [98] for more details.

4.3. Electroweak precision test – ΔS, ΔT and ΔU

Finally, we would like to comment on extra corrections from additional particles in G2HDM to the electroweak oblique observables. In additional to the SM particles, G2HDM contains the new $SU(2)_L$ doublet H_2, the $SU(2)_H$ gauge bosons of which Z' mixes with the SM Z, and the heavy $SU(2)_H$ fermions. Other scalars Φ_H and Δ_H are singlets under $SU(2)_L$ and hence are not relevant.

The heavy fermions, as $SU(2)_L$ singlets, will not contribute to electroweak corrections described by the oblique parameters, ΔS, ΔT and ΔU, as can be easily seen from the definition of the parameters [99]. Moreover, as demonstrated above the Z–Z' mixing is constrained by the electron $g-2$ bound to be less than 10^{-2} or so, implying contributions to the oblique parameters at the order of 10^{-4} or smaller. Finally as long as the mass splitting between $H_2^\pm$ and H_2^0 is small, corrections to ΔS, ΔT and ΔU will be suppressed [100]. All in all, this model can survive from the electroweak precision test.

5. Conclusion

In this work, we address a possible solution to the diphoton anomaly observed at the LHC based on the recent G2HDM model proposed by us. In the G2HDM, the two Higgs doublets H_1 and H_2 are embedded into a doublet under a non-abelian gauge symmetry $SU(2)_H$ and the resulting $SU(2)_H$ doublet is charged under an additional abelian group $U(1)_X$. To give masses to additional gauge bosons, we introduce a $SU(2)_H$ scalar triplet and a doublet (both are singlets under the SM gauge group). On the other hand, extra new heavy fermions are needed to have Yukawa couplings comply with the $SU(2)_H$ gauge symmetry. In other words, we have only *chiral* fermions, different from some of existing models where vector-like quarks and leptons are employed to explain the anomaly. In addition, constraints on new vector-like quarks and leptons because of mixing with SM fermions [101–103] do not apply here since our new fermions do not mix with the SM ones.

The new heavy fermions receive masses from the vev of the $SU(2)_H$ scalar doublet, that also provides masses to the additional gauge bosons. A physical component ϕ_2 inside the doublet can be produced radiatively via gluon fusion with the additional heavy colored fermions in loops and in turn radiatively decays into two photons with the heavy charged fermions involved. We have shown that in the limit of the universal fermion mass, in order to reproduce the anomaly, the vev of ϕ_2 ranges from 180 to 300 GeV with the new fermion mass of few hundred GeV. The desired total decay width of $\Gamma_{\phi_2} \simeq 0.06 m_{\phi_2}$, by having ϕ_2 decay into the new fermions, can be realized with $m_f \sim 360$ GeV and $v_\Phi \sim 250$ GeV. The favorable region could be further extended if the additional neutral fermions are allowed to have arbitrary masses.

The existence of $SU(2)_H$ gauge bosons can also explain the anomalous muon magnetic dipole moment. There are three radiative corrections to muon $g-2$: W' with μ^H, Z' with μ_R and the correction induced by the Z–Z' mixing. We have found out with $m_{\mu^H} = 360$ GeV and

$g_H \sim 7 \times 10^{-3}$, resulting in GeV or sub-GeV W' and Z' depending on the vevs of Φ_H and Δ_H, the muon anomaly Δa_μ of order 10^{-9} can be realized while the corresponding contributions to electron anomaly Δa_e are highly suppressed by the very small electron mass.

We conclude by pointing out that except for the diphoton anomaly, the LHC run-II data do not feature any significant deviation from the SM prediction. Our model can avoid overproducing other SM model particles through the same ϕ_2 exchange process since ϕ_2 couples only to the extra fermions at tree level in the limit of the vanishing h–ϕ_2 mixing. The heavy fermions from ϕ_2 decays, however, subsequently decay into SM particles plus the DM particles, that manifest as missing transverse energy. The resulting SM particle energy spectra depend on the mass difference between the new heavy fermions and DM, and the spectra could be very soft if the mass difference is small just like the compressed spectra in various supersymmetry models. Finally, for the zero h–ϕ_2 mixing, one can expect the $Z\gamma$ and ZZ signals with a similar order of magnitude as in the $\gamma\gamma$ anomaly through the same ϕ_2 exchange process.

Acknowledgments

WCH would like to thank Ivan Nišandžić for very useful discussion on parton distribution functions. This work is supported in part by the Ministry of Science and Technology (MOST) of Taiwan under grant number 104-2112-M-001-001-MY3 (TCY), DGF Grant No. PA 803/10-1 (WCH), and the World Premier International Research Center Initiative (WPI), MEXT, Japan (YST).

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