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See-saw composite Higgs model at the LHC: Linking naturalness to the 750 GeV diphoton resonance

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We explore the possibility of explaining the recent ~ 750 GeV excesses observed by ATLAS and CMS in the $\gamma\gamma$ spectrum in the context of a compelling theory of naturalness. The potential spin-zero resonance responsible for the excesses also requires the existence of new heavy charged states. We show that both such features are naturally realized in a see-saw composite Higgs model for electroweak symmetry breaking, where the new pseudo-Goldstone bosons are expected to be comparatively heavier than the Standard Model Higgs, and the new fermions have masses in the TeV range. If confirmed, the existence of this new resonance could be the first stone in the construction of a new theory of naturalness.

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I. INTRODUCTION

Very recently, both ATLAS [1] and CMS [2] Collaborations have observed a prominent excess in the diphoton spectrum around $m_{\gamma\gamma} \sim 750$ GeV, which could be the first signature of new physics beyond the Standard Model (SM) at the Large Hadron Collider (LHC). The most likely hypothesis is that of a spin-zero resonance produced via gluon fusion, which if confirmed would mean the discovery of a new scalar degree of freedom. Theories with scalar particles, including the Higgs, require a fine-tuning of parameters unless one introduces new symmetries or dynamics. With the unveiling of a new scalar resonance, substantially heavier than the Higgs, we are facing a new challenge to accommodate both scalars in a single natural setup.

Here we show how this can be achieved in the context of composite Higgs [3], based on the model proposed by some of the authors [4], with the Higgs and the new resonance realized as pseudo-Goldstone bosons of spontaneously broken global symmetries. This mechanism protects the scalar potential from the problematic UV sensitivity, and at the same time opens a door to address key open questions in the SM, such as the nature of dark matter (DM) and the origin of the matter-antimatter asymmetry in the Universe.

For these scenarios to be viable completions of the SM, the scalar potential must allow for electroweak symmetry breaking (EWSB), which requires a negative mass-squared term for the Higgs field. This is generically induced via fermionic loop contributions to the potential, coming dominantly from “top partner” states. In standard composite Higgs scenarios, in order to achieve natural EWSB together with a light Higgs mass, these new fermions cannot be much heavier than the electroweak (EW) scale $v = 246$ GeV, which creates a significant amount of tension as the current LHC limits push their mass towards the TeV scale.

In [4] an elegant solution to the above problem was proposed, in the form of a see-saw type of EWSB, based on

a sequential symmetry breaking pattern such as $SO(6) \rightarrow SO(5) \rightarrow SO(4)$, which gives rise to a pseudo-Goldstone doublet ϕ and singlet η from the first breaking, and another doublet θ from the second breaking. Due to the sequential pattern, the η , ϕ fields are expected to be significantly heavier than θ , since $SO(6)/SO(5)$ breaking interactions would generate a mass for η , ϕ but not θ . In this scenario, scalar potential terms of the form $\mu^2 \phi^\dagger \theta + \text{H.c.}$ would give rise to a mixing between the heavy and light scalar doublets, yielding after diagonalization a negative mass term for the light doublet eigenstate, which would trigger EWSB without the need of light top partners, those being now linked instead to the heavier pseudo-Goldstone scalars ϕ , η .

As a mass hierarchy between the Higgs and the supposed new scalar resonance around 750 GeV is precisely what one would expect in this scenario, it is compelling to investigate the possibility that the new resonance can be identified with either of the heavy scalars ϕ or η , possibility which we explore in this work.

II. A SEE-SAW COMPOSITE HIGGS MODEL

Let us now discuss the main features of our setup (for a more detailed discussion, see [4]). The model features a global $SO(6)$ symmetry that is spontaneously broken via $SO(6) \rightarrow SO(5) \rightarrow SO(4)$. This spontaneous breaking is assumed to be triggered by the condensation of some strongly interacting sector endowed with a global $SO(6)$ symmetry. The scales of the first and second breakings are denoted by F_1 and F_2 respectively. These scales correspond to the decay constants of the Goldstone bosons, and are generated dynamically via the strongly interacting dynamics. Although in the simplest scenario the breakings occur at the same scale (i.e. $F_1 = F_2$), we also allow for the possibility that the two sets of Goldstone bosons have *different* decay constants ($F_1 > F_2$). As discussed in [4], a

large hierarchy between the two scales $F_1 \gg F_2$ will however reintroduce some degree of tuning.

The first breaking $SO(6) \rightarrow SO(5)$ gives rise to five Goldstone bosons, an $SU(2)_L$ doublet ϕ and a singlet η [5], while the second breaking $SO(5) \rightarrow SO(4)$ gives rise to another doublet θ . The light doublet θ will eventually be associated with the SM-like Higgs doublet. The lightness of θ is guaranteed by the (approximate) $SO(5)$ symmetry, so that any corrections to its mass m_θ will be proportional to the amount of $SO(5)$ breaking. Large sources of $SO(6)/SO(5)$ breaking will (in the absence of any other explicit breaking) leave the potential for θ unaffected, while contributing to the potential for ϕ and η . Thus we consider scenarios in which $SO(6)/SO(5)$ is badly broken, since these naturally lead to a mass hierarchy between the $SO(6)/SO(5)$ Goldstones (ϕ and η) and the Higgs.

The explicit $SO(6)/SO(5)$ breaking could come from a variety of sources: a simple possibility is to consider a new set of elementary fermions ψ in a multiplet of $SO(5)$ coupled to fermionic strong sector operators \mathcal{O}_ψ via

$$\Delta\mathcal{L} = \alpha\bar{\psi}\mathcal{O}_\psi + \text{H.c.}, \quad (1)$$

such that loops of these fermions induce large contributions to $V(\phi, \eta)$ with a characteristic scale $\alpha F_1/4\pi$. Alternatively, the gauge group could be enlarged to include an extra set of gauge bosons A_μ coupling to ϕ and η , but not to θ

$$\Delta\mathcal{L} = \beta J_\mu^{\phi, \eta} A^\mu. \quad (2)$$

Here $J_\mu^{\phi, \eta}$ represents the current associated with the $SO(6)/SO(5)$ symmetry, which excites ϕ and η from the vacuum. In the presence of (2), contributions to the potential $V(\phi, \eta)$ would scale as $\beta F_1/4\pi$. The explicit breaking parameters α, β may be large without affecting the second Goldstone doublet θ , since the couplings in (1) and (2) preserve $SO(5)$.

As discussed above, the presence of sources of explicit $SO(6)/SO(5)$ breaking in the UV theory yield mass terms for the first set of Goldstone bosons

$$\mathcal{L}_{\phi, \eta} = m_\phi^2 \phi^2 + m_\eta^2 \eta^2, \quad (3)$$

as well as mixing terms between the various sets of Goldstones

$$\begin{aligned} \mathcal{L}_{\text{mix}} = & A_1 F_2 \phi \cdot \theta \frac{s_\theta}{|\theta|} + A_2 F_2 \eta c_\theta + B_1 F_2^2 (\phi \cdot \theta)^2 \frac{s_\theta^2}{|\theta|^2} \\ & + B_2 F_2^2 \eta^2 c_\theta^2 + 2B_3 F_2^2 \eta \phi \cdot \theta \frac{s_\theta c_\theta}{|\theta|}, \end{aligned} \quad (4)$$

where the Goldstone doublets ϕ, θ have been expressed as vectors $\phi = (\phi^1 \phi^2 \phi^3 \phi^4)^T$ and $\theta = (\theta^1 \theta^2 \theta^3 \theta^4)^T$ of $SO(4) \simeq SU(2)_L \times SU(2)_R$. The parameters A_i and B_i have mass dimension $[A] = 2$, $[B] = 0$, and

$$s_\theta = \sin \frac{|\theta|}{F_2}, \quad c_\theta = \cos \frac{|\theta|}{F_2}. \quad (5)$$

In the exact $SO(5)$ limit we must have $A_1 = A_2$ and $B_1 = B_2 = B_3$ in (4), as well as $m_\phi = m_\eta$ in (3). In this limit (3) and (4) yield the most general effective Lagrangian for ϕ, θ and η up to quartic order in the fields and invariant under $SO(5)$. The invariance under $SO(5)$ can be made manifest by grouping ϕ and η together into a vector of $SO(5)$: $(\phi^1 \phi^2 \phi^3 \phi^4 \eta)$, while θ is parametrized by a nonlinear sigma field Σ

$$\begin{aligned} \Sigma(\theta) = & \exp(i\sqrt{2}X^a \theta^a / F_2)(0, 0, 0, 0, 1)^T \\ = & \frac{\sin(|\theta|/F_2)}{|\theta|/F_2} (\theta^1, \theta^2, \theta^3, \theta^4, |\theta| \cot(|\theta|/F_2))^T, \end{aligned} \quad (6)$$

where X^a are the spontaneously broken $SO(5)/SO(4)$ generators. Σ parametrizes the fluctuations of θ around the $SO(5)$ breaking vacuum, and transforms as a vector of $SO(5)$. These transformation properties guarantee the $SO(5)$ invariance of (4) in the specified limit.

Assuming all the $SO(6)/SO(5)$ breaking effects to have a common origin, the dimensionful parameters A_1 and m_ϕ^2 [respectively equal to A_2 and m_η^2 in the $SO(5)$ invariant limit] will be of similar size. As an example, if the breaking is due to an interaction such as (1), both would be of order $(\alpha F_1/4\pi)^2$. In turn, coupling the strong sector to the SM will induce an explicit breaking of $SO(5)$. Thus, deviations from the $SO(5)$ invariant limit in (4) are expected comparable in size to the loop induced mass of the light doublet. If we define $\delta m^2 = m_\phi^2 - m_\eta^2$ and $\delta A = A_1 - A_2$, then $|\delta m^2| \approx |\delta A| \approx m_\phi^2$. In summary, all contributions to the light Higgs mass are of order $\sqrt{\delta A}, \sqrt{\delta m^2}$, being induced by $SO(5)$ breaking effects, while the masses of ϕ and η remain tied to the scale of $SO(6)$ breaking and can therefore be parametrically larger. The model thus features a spectrum of scalar particles that exhibits a natural hierarchy of scales.

Looking at (4), we first note that for $A_2 \neq 0$, the singlet field η develops a vacuum expectation value (vev), $\eta \rightarrow \langle \eta \rangle + \eta$, with

$$\langle \eta \rangle = -\frac{A_2 F_2}{2(m_\eta^2 + B_2 F_2^2)}, \quad (7)$$

and bearing in mind that $A_2 \sim m_\eta^2$, we have $|\langle \eta \rangle| \lesssim F_2/2$. At the same time, the term proportional to A_1 in (4) induces a mixing between the two doublets ϕ, θ . The mass matrix reads

$$\begin{pmatrix} m_\phi^2 & \mu^2 \\ \mu^2 & m_\theta^2 - \frac{A_2}{2F_2} \langle \eta \rangle - B_2 \langle \eta \rangle^2 \end{pmatrix} \quad (8)$$

with $\mu^2 \equiv A_1/2 + B_3 F_2 \langle \eta \rangle$. The mixing yields two (doublet) eigenstates H and h , the latter being the light SM-like Higgs, which obtains a negative mass squared for $\mu^4 > m_\phi^2 \times (m_\theta^2 - A_2 \langle \eta \rangle / (2F_2) - B_2 \langle \eta \rangle^2)$. The occurrence of such a negative mass-squared term from the mixing of the two Goldstone $SU(2)_L$ doublets, associated with the sequential global symmetry breaking pattern, is key in this framework, yielding viable EWSB *à la* see-saw (see [6,7] for a similar realization of EWSB in other contexts). The rotation to the doublet mass eigenbasis is given by $\phi = c_\alpha \mathbf{H} - s_\alpha \mathbf{h}$, $\theta = s_\alpha \mathbf{H} + c_\alpha \mathbf{h}$, with $c_\alpha \equiv \cos \alpha$, $s_\alpha \equiv \sin \alpha$ and the rotation angle given by

$$\tan 2\alpha = \frac{A_1 + 2B_3 F_2 \langle \eta \rangle}{m_\phi^2 + \frac{A_2}{2F_2} \langle \eta \rangle + B_2 \langle \eta \rangle^2 - m_\theta^2}. \quad (9)$$

Expanding the scalar potential (4) we find that the relevant terms involving η , H and h (up to $D = 4$) are

$$\begin{aligned} V(H, h, \eta) = & -\mu_h^2 h^\dagger h + \mu_H^2 H^\dagger H + (m_\eta^2 + B_2 F_2^2) \eta^2 \\ & - \left[\left(\frac{A_2}{2F_2} + 2B_2 \langle \eta \rangle \right) c_\alpha^2 + B_3 F_2 s_{2\alpha} \right] \eta h^\dagger h \\ & - \left[\left(\frac{A_2}{2F_2} + 2B_2 \langle \eta \rangle \right) s_\alpha^2 - B_3 F_2 s_{2\alpha} \right] \eta H^\dagger H \\ & - \left[\left(\frac{A_2}{2F_2} + 2B_2 \langle \eta \rangle \right) s_\alpha c_\alpha - 2B_3 F_2 c_{2\alpha} \right] \eta H^\dagger h \\ & + B_2 c_\alpha^2 \eta^2 h^\dagger h - B_2 s_\alpha^2 \eta^2 H^\dagger H \\ & - B_2 s_{2\alpha} \eta^2 H^\dagger h + \text{H.c.} \end{aligned} \quad (10)$$

with $-\mu_h^2$ and μ_H^2 the resulting squared-mass terms after the diagonalization of (8).

III. COUPLINGS OF H AND η TO $\gamma\gamma$

The couplings of the new heavy scalar states H and η to $\gamma\gamma$ occur via loops of the heavy fermions Ψ responsible for the explicit $SO(6)/SO(5)$ breaking. These generically transform both under $SU(3)_C$ and $U(1)_Y$. We note that in order to have consistent hypercharge assignments we need to extend the global symmetry to include an extra $U(1)_X$. The SM gauge group is embedded in $SO(6) \times U(1)_X$ in such a way that hypercharge is realized as $Y = X + T^{3R}$, where T^{3R} is the third component of the $SU(2)_R$ subgroup of the custodial $SO(4) \sim SU(2)_L \times SU(2)_R$.

Transforming under $SU(3)_C$ and $U(1)_Y$, the heavy fermions may be responsible both for the production of these scalars at the LHC in gluon fusion $pp(gg) \rightarrow H, \eta$, and their subsequent decay into $\gamma\gamma$. The effective couplings of $H, \eta = \phi$ to SM gauge bosons are given by

$$\mathcal{L}_G = -\frac{c_1 \alpha_1}{4} \phi B_{\mu\nu} B^{\mu\nu} - \frac{c_2 \alpha_2}{4} \phi W_{\mu\nu}^i W^{i\mu\nu} - \frac{c_3 \alpha_3}{4} \phi G_{\mu\nu}^a G^{a\mu\nu} \quad (11)$$

with $i = 1, 2, 3$, $a = 1, \dots, 8$ and $\alpha_{1,2}$ being respectively $g^2/(4\pi)$ and $g^2/(4\pi)$. We note that (11) assumes η to be a CP -even state, whereas if η is a CP -odd particle, we need to substitute one of the field strengths by a dual: $B_{\mu\nu} B^{\mu\nu} \rightarrow B_{\mu\nu} \tilde{B}^{\mu\nu}$, $W_{\mu\nu}^i W^{i\mu\nu} \rightarrow W_{\mu\nu}^i \tilde{W}^{i\mu\nu}$ and $G_{\mu\nu}^a G^{a\mu\nu} \rightarrow G_{\mu\nu}^a \tilde{G}^{a\mu\nu}$. We can reexpress the interactions in (11) in terms of the physical SM gauge bosons as

$$\begin{aligned} \mathcal{L}_G = & -\frac{g_{\gamma\gamma}^\phi}{4} \phi F_{\mu\nu} F^{\mu\nu} - \frac{g_{Z\gamma}^\phi}{4} \phi F_{\mu\nu} Z^{\mu\nu} - \frac{g_{ZZ}^\phi}{4} \phi Z_{\mu\nu} Z^{\mu\nu} \\ & - \frac{g_{W\gamma}^\phi}{4} \phi W_{\mu\nu} W^{\mu\nu} - \frac{g_G^\phi}{4} \phi G_{\mu\nu}^a G^{a\mu\nu} \end{aligned} \quad (12)$$

with $g_{\gamma\gamma}^\phi = c_1 \alpha_1 c_W^2 + c_2 \alpha_2 s_W^2$, $g_{Z\gamma}^\phi = (c_1 \alpha_1 - c_2 \alpha_2) s_{2W}$, $g_{ZZ}^\phi = c_1 \alpha_1 s_W^2 + c_2 \alpha_2 c_W^2$, $g_{W\gamma}^\phi = 2c_2 \alpha_2$, $g_G^\phi = c_3 \alpha_3$.

The relation between the effective operators in (12) and the heavy fermions Ψ depends on the specific fermion representation under the global symmetry group, their transformation properties under the SM gauge symmetries and whether the scalars ϕ acquire a vev. Among our fermionic bound states, all colored fermions will participate in the coupling of the scalar resonance to gluons in a universal fashion, and in the following we will denote their number as N_3 , whereas the number of fermions contributing to the EW couplings will be denoted by N_{EW} .

Focusing on fermions in a spinorial $\mathbf{4}$ of $SO(5)$ [8] (which we denote as Ψ_4), and denoting by X their $U(1)_X$ charge, the effective couplings of $\phi = \eta$ to EW gauge bosons are given in Table I. We consider the fermions Ψ_4 to have a common mass M_Ψ , linked to the strong dynamics responsible for the breaking of $SO(6) \rightarrow SO(5)$, and thus $M_\Psi \sim F_1 \gg v$ naturally. We note that for a nonvanishing $\langle \eta \rangle$, the heavy fermions get a correction to their mass term

$$y_\eta \langle \eta \rangle \bar{\Psi}_4 \gamma_5 \Psi_4 = y_\eta \langle \eta \rangle (\bar{\psi}^1 \psi^1 + \bar{\psi}^2 \psi^2 - \bar{\psi}^3 \psi^3 - \bar{\psi}^4 \psi^4), \quad (13)$$

which is anyway subdominant since $y_\eta |\langle \eta \rangle| / M_\Psi \ll 1$ (recall $|\langle \eta \rangle| < F_2/2$, $M_\Psi \sim F_1$ and $F_1 > F_2$).

We can also consider the coupling of $\phi = H$ to two photons. We first note that the eigenstate ϕ_4 coupling to the heavy fermions does it as

TABLE I. Couplings of η to EW gauge bosons.

	$c_1 / (\frac{y_\eta N_{EW}}{M_\Psi})$	$c_2 / (\frac{y_\eta N_{EW}}{M_\Psi})$	$g_{\gamma\gamma}^\eta / (\frac{\alpha_{EM} y_\eta N_{EW}}{M_\Psi})$
$\langle \eta \rangle = 0$	$-\frac{1}{2}$	$\frac{1}{2}$	0
$\langle \eta \rangle \neq 0$	$-\frac{1}{2} - (\frac{1}{2} + 4X^2) \frac{y_\eta \langle \eta \rangle}{M_\Psi}$	$\frac{1}{2} - \frac{y_\eta \langle \eta \rangle}{2M_\Psi}$	$-(1 + 4X^2) \frac{y_\eta \langle \eta \rangle}{M_\Psi}$

TABLE II. Ratios r_{XY} for H and η .

	r_{ZZ}	$r_{Z\gamma}$	r_{WW}
η	$\frac{2.718}{(1+4X^2)^2} r_\eta$	$\frac{1.9}{(1+4X^2)^2} r_\eta$	$\frac{21.11}{(1+4X^2)^2} r_\eta$
H	$\frac{3.682+4.356X^2+1.289X^4}{(1+4X^2)^2}$	$\frac{5.917+20.77X^2+18.24X^4}{(1+4X^2)^2}$	$\frac{21.11}{(1+4X^2)^2}$

$$y_\phi \phi_4 \bar{\Psi}_4 \gamma_4 \Psi_4 = y_\phi \phi_4 (\bar{\psi}^1 \psi^3 + \bar{\psi}^2 \psi^4) + \text{H.c.} \quad (14)$$

The mixing between ϕ and θ yields a correction to the heavy fermion mass term after EWSB

$$y_\phi s_\alpha v (\bar{\psi}^1 \psi^3 + \bar{\psi}^2 \psi^4) + \text{H.c.}, \quad (15)$$

where $v = \langle h \rangle$ is the Higgs vev. In this case

$$c_1 = (1/2 + 4X^2) y_\phi^2 \frac{s_\alpha v}{M_\Psi^2}, \quad c_2 = \frac{y_\phi^2 s_\alpha v}{2 M_\Psi^2}, \quad (16)$$

which yield a coupling of H to photons given by

$$g_{\gamma\gamma}^H = -\frac{N_{\text{EW}} y_\phi^2 s_\alpha v}{M_\Psi^2} (1 + 4X^2) \alpha_{\text{EM}}.$$

From (12), the partial decay width of either H or η into photons is given by

$$\Gamma(\varphi \rightarrow \gamma\gamma) = \frac{(g_{\gamma\gamma}^\varphi)^2}{64\pi} m_\varphi^3. \quad (17)$$

The relation between the branching ratio of either H or η into photons and into other vector bosons, namely $r_{XY} = \Gamma(\varphi \rightarrow XY)/\Gamma(\varphi \rightarrow \gamma\gamma)$ is shown in Table II, with $r_\eta \equiv M_\Psi^2/(y_\eta \langle \eta \rangle)^2$.

For a canonical choice $X = \pm 1/2$, the ratios for the heavy Higgs H are $r_{ZZ} = 1.2$, $r_{Z\gamma} = 3.1$ and $r_{WW} = 5.3$, whereas for the η particle they are a function of the fermion masses and couplings, namely $(r_{ZZ}, r_{Z\gamma}, r_{WW}) = r_\eta \times (0.7, 0.5, 5.3)$.

IV. DIPHOTON SIGNATURES AT THE LHC

Using the results from the previous section, we now analyze the possibility that either H or η in our framework correspond to the potential diphoton resonance observed by both ATLAS and CMS around $m_\varphi \sim 750$ GeV. We first note that in order for any new scalar φ to have a sizeable branching fraction into $\gamma\gamma$, its tree-level decays into other SM particles should be absent or heavily suppressed. Then, for $\varphi = \eta$, the term $\eta h^\dagger h$ in (10) poses a potentially important obstacle towards achieving a sizeable $\text{Br}(\eta \rightarrow \gamma\gamma)$. The partial width $\Gamma(\eta \rightarrow hh)$ is given by

$$\Gamma(\eta \rightarrow hh) = \frac{\kappa_{\eta hh}^2}{8\pi m_\eta} \sqrt{1 - \frac{4m_h^2}{m_\eta^2}} \quad (18)$$

with

$$\kappa_{\eta hh} = -\left[\left(\frac{A_2}{2F_2} + 2B_2 \langle \eta \rangle \right) c_\alpha^2 + B_3 F_2 s_{2\alpha} \right], \quad (19)$$

such that the relation $\Gamma(\eta \rightarrow hh) \lesssim \Gamma(\eta \rightarrow \gamma\gamma)$ would lead to

$$\kappa_{\eta hh} \lesssim \frac{Q_\Psi^4 y_\eta^2 \alpha_{\text{EM}}^2 m_\eta^2}{4\pi^2 M_\Psi} \quad (20)$$

with $Q_\Psi = \sqrt{1 + 4X^2}$ being the electric charge of the heavy fermions running in the loop which mediates $\eta \rightarrow \gamma\gamma$. After EWSB $\kappa_{\eta hh} \neq 0$ also gives rise to singlet-doublet mixing, such that the singletlike mass eigenstate inherits a small amount of the Higgs couplings to SM particles. While the value of this mixing β is constrained by a combination of LHC measurements of Higgs signal strengths and EW precision observables to $s_\beta < 0.32$ at 95% C.L. for $m_\eta \sim 750$ GeV [9], admixtures below this value may still yield $\Gamma(\eta \rightarrow WW, ZZ, i\bar{i}) \gg \Gamma(\eta \rightarrow \gamma\gamma)$.

We emphasize here that $\text{Br}(\eta \rightarrow \gamma\gamma) \ll 1$ does not necessarily rule out the possibility of accounting for the diphoton excess (on the contrary, if the ATLAS favored diphoton width $\Gamma_{\gamma\gamma} \sim 45$ GeV [1] is confirmed, this will generically imply $\text{Br}(\varphi \rightarrow \gamma\gamma) \ll 1$). However, given the diphoton cross section compatible with the excess, whose best-fit value is given by [10]

$$\sigma(pp \rightarrow \varphi \rightarrow \gamma\gamma) = 6.2 \pm 1.0 \text{ fb}, \quad (21)$$

current limits from LHC Run 1 searches in other final states may yield stringent limits on the value of X and r_η . In particular, $r_{ZZ} \lesssim 13$, $r_{Z\gamma} \lesssim 7$, $r_{WW} \lesssim 45$, $r_{hh} \lesssim 41$ at 95% C.L. from LHC searches at 8 TeV (see e.g. [10]), such that for $X = \pm 1/2$, $r_\eta \lesssim 10$ is required [11] to satisfy these bounds. LHC Run 2 may nevertheless be able to explore decays of the resonance beyond $\gamma\gamma$, into other states such as WW , ZZ and $Z\gamma$.

For $\varphi = H$, since H and h do not mix (by construction), H does not have *a priori* any dangerous tree-level decays into SM particles. Moreover, in this case the constraints from Run 1 on r_{XY} are automatically satisfied for $X = \pm 1/2$ [and in general for $X \sim \mathcal{O}(1)$]. The cross section $\sigma(pp \rightarrow H \rightarrow \gamma\gamma)$ needed to accommodate the diphoton excess, its best-fit value given by (21), can be related to the fermion parameters responsible for the production and decay. The production cross section $\sigma_{XS}(gg \rightarrow H)$ is only dependent on g_G^H in (12), and its value at 13 TeV LHC is (see e.g. [10])

$$\sigma_{XS}(gg \rightarrow H) \simeq \left(\frac{g_G^H}{\text{TeV}^{-1}} \right)^2 \times 100 \text{ pb}. \quad (22)$$

The effective coupling g_G^H can be expressed in terms of $g_{\gamma\gamma}^H$ as $g_G^H = g_{\gamma\gamma}^H \times \alpha_3 / (6\alpha_{\text{EM}} Q_\Psi^2) \gg g_{\gamma\gamma}^H$. This allows us to write

$\sigma(pp \rightarrow H \rightarrow \gamma\gamma)$ solely as a function of $g_{\gamma\gamma}^H$ (assuming the dominant decay mode of H is into gluons, which is the case in our scenario)

$$\sigma(pp \rightarrow H \rightarrow \gamma\gamma) \simeq \left(\frac{g_{\gamma\gamma}^H}{\text{TeV}^{-1}} \right)^2 \times 13 \text{ pb.} \quad (23)$$

Combining (21) and (23) results in a preferred value $(g_{\gamma\gamma}^H)^{-1} \sim 45 \text{ TeV}$. Using the results from Sec. III, assuming $X = \pm 1/2$ and $y_\phi \sim \mathcal{O}(1)$, this translates into

$$(M_\Psi/v)^2 \simeq 2N_{\text{EW}}. \quad (24)$$

As we expect $M_\Psi \sim F_1 \sim \text{TeV}$, this leads to an estimate on the number of degrees of freedom contributing to the diphoton coupling, namely $N_{\text{EW}} \simeq \mathcal{O}(5 - 10)$, which incidentally matches the expectation from fermions transforming under a low representation of $SO(5)$.

V. ASTROPHYSICAL AND COSMOLOGICAL CONSEQUENCES

There are a number of possible cosmological and astrophysical consequences of this scenario which deserve a more detailed study, and which we discuss below.

Let us discuss first the implications for DM. The neutral heavy fermions in our model can play the role of DM, with the resonance φ playing the role of a DM mediator. Similar scenarios have been discussed in the literature in the context of radion/dilaton and axion mediators, for the CP -even [12] and CP -odd [13] cases. In our case, the DM mediator should be a CP -odd η , since then the annihilation cross section is s -wave unsuppressed. In this context a simple choice is then $X = \pm 1/2$ for one fermion multiplet, which leads to two neutral fermions. Among these, the lightest one will be DM, with a small splitting with the next state of order $(y_\phi v s_\alpha / M_\Psi)^2$ or $(y_\eta \langle \eta \rangle / M_\Psi)^2$. This will lead to a model similar to inelastic DM [14] or pseudo-Dirac DM [15], with coannihilations playing an important role. The main annihilation process would be to gluons, as φ decays predominantly to gluons. The relic abundance is then proportional to the combination $4\pi^3 M_\Psi^2 / (y_\eta \alpha_s)^2$. Values in the range $M_\Psi \sim \text{TeV}$ and $y_\eta \lesssim \mathcal{O}(1)$ lead to a relic abundance in agreement with Planck [16].

DM in this scenario would also produce γ rays via the coupling of φ to $\gamma\gamma$ and $Z\gamma$. Understanding the correlations

of the 750 GeV signal with possible lines in the spectrum measured by Fermi-LAT and HESS [17] could lead to a selection of fermionic representations in this model.

Finally, there is a tantalizing correlation between DM and baryogenesis in this model. As we mentioned before, efficient annihilation requires a pseudoscalar η mediator, and the generation of the diphoton signal implies η would get a vev, hence breaking spontaneously CP . Additional fermionic states, new scalars and CP violation are excellent starting points to explore EW baryogenesis [18] in this model.

VI. SUMMARY

In this paper we have presented an explanation of the diphoton signal seen by ATLAS and CMS in terms of a fully natural composite Higgs model. The model features a new spectrum of composite scalars, with masses of order TeV. We find that these new states can decay via loops of vectorlike heavy fermions and reproduce the observed diphoton excess. The mass hierarchy between the Higgs-like doublet and the new scalars is a crucial and natural feature of the see-saw composite Higgs model, and thus the new states are completely natural components of the model.

We have also identified a potential dark matter candidate: with a suitable $U(1)_X$ charge assignment the vectorlike fermions can form neutral states that will behave as inelastic/pseudo-Dirac DM. For natural values of the model parameters we find that the model leads to successful relic abundance.

Arriving at a satisfying solution to the hierarchy problem without resorting to fine-tuning is a long standing challenge. Most potential solutions to the problem lead to us to expect new resonances around the TeV scale. If the recent diphoton signal is the first such observation, we believe the model we have presented succeeds in explaining the data in a coherent, and most importantly natural, fashion.

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