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Dynamical seesaw mechanism for Dirac neutrinos

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A B S T R A C T

So far we have not been able to establish that, as theoretically expected, neutrinos are their own anti-particles. Here we propose a dynamical way to account for the Dirac nature of neutrinos and the smallness of their mass in terms of a new variant of the seesaw paradigm in which the energy scale of neutrino mass generation could be accessible to the current LHC experiments.

1. Introduction

Although neutrinos are not proven experimentally to be Dirac or Majorana fermions, the leading theoretical expectation is that they are Majorana. Indeed there is a widespread paradigm that ascribes the smallness of neutrino masses relative to the other Standard Model fermion masses to their charge neutrality. This is naturally incorporated by assuming that neutrinos acquire Majorana masses from a lepton number violating operator, such as Weinberg's dimension five operator or similar higher order ones. Indeed conventional type-I [1–5] or type-II [4–6] formulations of the seesaw mechanism lead typically to Majorana neutrinos, irrespective of whether the seesaw is realized at high or at low mass scale, in the spirit of the models considered in [7]. Until the observation of neutrinoless double beta decay [8] becomes unambiguously confirmed [9] the possibility remains that neutrinos can be Dirac particles after all.

The theoretical challenge to account for this possibility is then twofold: i) to predict Dirac neutrinos, and ii), to understand dynamically their small mass. Regarding the first we need to use extra symmetries beyond $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ gauge symmetry, otherwise massive neutrinos are generally expected to be Majorana particles [4]. To this end, within the standard $SU(3)_c \otimes SU(2)_L \otimes U(1)_Y$ electroweak gauge structure, one may impose a conserved lepton number, so as to obtain Dirac neutrinos. If the lepton number assignment is non-standard one may obtain naturally light Dirac neutrinos from calculable radiative corrections [10]. Likewise, one may consider schemes based on flavor symmetries, as suggested in [11]. Unfortunately in the simplest realization of this idea the smallness of neutrino mass is put in by hand. Another approach would be to appeal to the existence of extra dimensions [12,13].

Alternatively one may extend the gauge group itself, so as to (at least partially) include the lepton number symmetry, for example, by using the extended $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ gauge structure [14] which predicts the number of fermion generations to match the number of colors. Although modern $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ schemes have Majorana neutrinos, with either radiatively induced [15,16] or seesaw-type masses [17,18], the simplest original formulation predicts Dirac neutrino masses [19]. Apart from not being able to account for current oscillation data, that original formulation did not address the question of how to account for the observed smallness of neutrino mass.

In this letter we focus on the possibility of having naturally light Dirac neutrinos with masses induced *a la seesaw*. We adopt the $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ gauge structure because of its unique features, with respect to other electroweak extensions based, for example, on left-right symmetry. Indeed, our new variant of the seesaw mechanism for Dirac neutrinos makes use of the peculiar features of the $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ based models. The resulting scheme provides a Dirac seesaw alternative to the approaches considered in [15,16] and [17,18].

2. The model

Our starting point is a variant of the $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ gauge framework with the anomaly-free matter content given in Table 1. Notice that one has the same set of "gauge-charged" fields as in [15], for example the left-handed leptons transform as

$$\psi_L^\ell = \begin{pmatrix} \ell^- \\ \nu_\ell \\ N_\ell \end{pmatrix}_L , \qquad (1)$$

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Table 1
Matter content of the model, where $\hat{u}_R \equiv (u_R, c_R, t_R, U_R)$ and $\hat{d}_R \equiv (d_R, s_R, b_R, D_R, D'_R)$.

	ψ_L^ℓ	ℓ_R	$S_R^\ell, \tilde{S}_R^\ell$	$Q_L^{1,2}$	Q_L^3	$\hat{u}_R$	$\hat{d}_R$	ϕ_0	ϕ_1	ϕ_2
$SU(3)_c$	**1**	**1**	**1**	**3**	**3**	**3**	**3**	**1**	**1**	**1**
$SU(3)_L$	$\mathbf{3}^*$	**1**	**1**	**3**	$\mathbf{3}^*$	**1**	**1**	$\mathbf{3}^*$	$\mathbf{3}^*$	$\mathbf{3}^*$
$U(1)_X$	$-\frac{1}{3}$	-1	0	0	$+\frac{1}{3}$	$+\frac{2}{3}$	$-\frac{1}{3}$	$+\frac{2}{3}$	$-\frac{1}{3}$	$-\frac{1}{3}$
$\mathcal{L}$	$-\frac{1}{3}$	-1	1	$-\frac{2}{3}$	$+\frac{2}{3}$	0	0	$+\frac{2}{3}$	$-\frac{4}{3}$	$-\frac{4}{3}$
$\mathbb{Z}_3^{\mathrm{aux}}$	ω	ω	ω	ω^2	ω^2	ω^2	ω^2	1	1	1

with $\ell = 1, 2, 3 \equiv e, \mu, \tau$. In contrast to what is assumed in [15], however, we add more gauge singlet leptons and we change the transformation properties of two of the scalar multiplets under $\mathcal{L}$, the ungauged piece of the lepton number symmetry. Indeed, in this model the electric charge can be written in terms of the $U(1)_X$ generator X and the diagonal generators of the $SU(3)_L$, whereas lepton number has a gauge component as well as a complementary global one:

$$Q = T_3 + \frac{1}{\sqrt{3}} T_8 + X \,, \tag{2}$$

$$L = \frac{4}{\sqrt{3}} T_8 + \mathcal{L} \,. \tag{3}$$

Let us now explain the new features of the model in relation to previous variants of the $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ electroweak model.

Change in the matter sector: In addition to the two-component neutral fermions N_L needed to fill up the $SU(3)_L$ anti-triplets, in our new model we introduce *two* sequential sets of $\mathcal{L}$-carrying gauge singlet leptons denoted as S_R, $\tilde{S}_R$.

Change in the scalar sector: As before, three scalar anti-triplets $\phi_0 \sim (\mathbf{3}^*, +2/3)$ and $\phi_{1,2} \sim (\mathbf{3}^*, -1/3)$ are responsible for the spontaneous breakdown of the extended electroweak $SU(3)_L$ gauge symmetry. However, in contrast to the formulation presented in [15], in the present framework the scalar triplets ϕ_1 and ϕ_2 have the same $\mathcal{L}$ charge. Following the notation of [19] the most general pattern for the vacuum expectation values (VEVs) of the fields is

$$\langle\phi_0\rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} k_0 \\ 0 \\ 0 \end{pmatrix}, \quad \langle\phi_1\rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ k_1 \\ n_1 \end{pmatrix}, \quad \langle\phi_2\rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ k_2 \\ n_2 \end{pmatrix}, \tag{4}$$

where the isosinglet VEVs n_1 and n_2 characterize the $SU(3)_L$ breaking scale. Correspondingly k_0, k_1 and k_2 are the VEVs of the $SU(2)_L \subset SU(3)_L$ doublets, so we expect that $k_0, k_1, k_2 \ll n_1, n_2$.

In the fermion sector, the relevant Yukawa terms invariant under the gauge symmetry and the auxiliary $\mathbb{Z}_3^{\mathrm{aux}}$ are

$$\begin{aligned} -\mathcal{L}_{\mathrm{f}} = {} & y^\ell \bar{\psi}_L l_R \phi_0 + y_1 \bar{\psi}_L S_R \phi_1 + \tilde{y}_1 \bar{\psi}_L S_R \phi_2 \\ & + y_2 \bar{\psi}_L \tilde{S}_R \phi_1 + \tilde{y}_2 \bar{\psi}_L \tilde{S}_R \phi_2 \\ & + y^u \bar{Q}_L^{1,2} \hat{u}_R \phi_0^* + \hat{y}^u \bar{Q}_L^3 \hat{u}_R \phi_1 + \tilde{y}^u \bar{Q}_L^3 \hat{u}_R \phi_2 \\ & + y^d \bar{Q}_L^3 \hat{d}_R \phi_0 + \hat{y}^d \bar{Q}_L^{1,2} \hat{d}_R \phi_1^* + \tilde{y}^d \bar{Q}_L^{1,2} \hat{d}_R \phi_2^* + \text{h.c.}, \end{aligned} \tag{5}$$

where contraction of the flavor indices is implicitly assumed. An important feature of the model is the fact that the $\mathcal{L}$ symmetry is preserved in the lepton sector. The role of the discrete symmetry $\mathbb{Z}_3^{\mathrm{aux}}$ is to forbid the gauge invariant $\psi_L^T C^{-1} \psi_L \phi_0$ term in order to ensure the seesaw suppression of the neutrino mass.

Concerning the symmetry breaking sector, the most general CP conserving scalar potential compatible with the $SU(3)_c \otimes SU(3)_L \otimes U(1)_X$ gauge symmetry as well as $\mathcal{L}$ invariance is given by

$$\begin{aligned} V_{\mathcal{L}} = {} & \sum_i \left(\mu_i^2 |\phi_i|^2 + \lambda_i |\phi_i|^4 \right) + \sum_{i \neq j} \lambda_{ij} |\phi_i|^2 |\phi_j|^2 \\ & + \sum_{i \neq j} \tilde{\lambda}_{ij} (\phi_i^\dagger \phi_j)(\phi_j^\dagger \phi_i) \\ & + \mu_s^2 \left(\phi_1^\dagger \phi_2 + \text{h.c.} \right) + \sum_i \lambda'_i \left(\phi_1^\dagger \phi_2 + \text{h.c.} \right) |\phi_i|^2 \\ & + \lambda \left[(\phi_1^\dagger \phi_2)(\phi_1^\dagger \phi_2) + \text{h.c.} \right] + \tilde{\lambda} \left[(\phi_1^\dagger \phi_0)(\phi_0^\dagger \phi_2) + \text{h.c.} \right], \end{aligned} \tag{6}$$

with $i = 0, 1, 2$. The key ingredient of our present construction is the inclusion of a $\mathcal{L}$ violating piece in the scalar potential:

$$V_{\not{\mathcal{L}}} = f \left(\phi_0 \phi_1 \phi_2 + \text{h.c.} \right), \tag{7}$$

where the term $\phi_0 \phi_1 \phi_2$ is understood as the fully antisymmetric product of the scalar field components. The mass dimension one parameter f is expected to be small, in the sense that the $\mathcal{L}$ symmetry is restored in the scalar sector in the limit $f \to 0$.[1] This symmetry can be interpreted as a chiral symmetry which will control the massiveness of neutrinos, as we show in the next section.

For simplicity, all VEVs and Yukawa coefficients are assumed to be real in the following analysis. From the minimization conditions of the total potential $V = V_{\mathcal{L}} + V_{\not{\mathcal{L}}}$: $\partial V/\partial k_0 = \partial V/\partial n_1 = \partial V/\partial n_2 = \partial V/\partial k_2 = 0$, the dimensionful parameters $\mu_{0,1,2}$ and f are determined as

$$\begin{aligned} \mu_0^2 = {} & -\frac{1}{2} \Bigg\{ 2\lambda_0 k_0^2 + \lambda_{01} \left(k_1^2 + n_1^2 \right) + \lambda_{02} \left(k_2^2 + n_2^2 \right) \\ & + \frac{\lambda'_0 \left[k_1^2 \left(2k_2^2 + n_2^2 \right) + 2k_1 k_2 n_1 n_2 + n_1^2 \left(k_2^2 + 2n_2^2 \right) \right]}{k_1 k_2 + n_1 n_2} \Bigg\} \\ & - \frac{(k_2 n_1 - k_1 n_2)^2}{2k_0^2} \Bigg[\tilde{\lambda}_{12} + 2\lambda \\ & + \frac{2\mu_s^2 + \lambda'_1 \left(k_1^2 + n_1^2 \right) + \lambda'_2 \left(k_2^2 + n_2^2 \right)}{k_1 k_2 + n_1 n_2} \Bigg], \\ \mu_1^2 = {} & -\frac{1}{2} \Bigg\{ \lambda_{01} k_0^2 + 2\lambda_1 \left(k_1^2 + n_1^2 \right) + (2\lambda + \lambda_{12} + \tilde{\lambda}_{12}) \left(k_2^2 + n_2^2 \right) \\ & + \frac{\left(k_2^2 + n_2^2 \right) \left[\lambda'_0 k_0^2 + \lambda'_2 \left(k_2^2 + n_2^2 \right) + 2\mu_s^2 \right]}{k_1 k_2 + n_1 n_2} \Bigg\} \\ & - \frac{\lambda'_1 \left[k_1^2 \left(3k_2^2 + n_2^2 \right) + 4k_1 k_2 n_1 n_2 + n_1^2 \left(k_2^2 + 3n_2^2 \right) \right]}{2(k_1 k_2 + n_1 n_2)}, \\ \mu_2^2 = {} & -\frac{1}{2} \Bigg\{ \lambda_{02} k_0^2 + 2\lambda_2 \left(k_2^2 + n_2^2 \right) + (2\lambda + \lambda_{12} + \tilde{\lambda}_{12}) \left(k_1^2 + n_1^2 \right) \\ & + \frac{\left(k_1^2 + n_1^2 \right) \left[\lambda'_0 k_0^2 + \lambda'_1 \left(k_1^2 + n_1^2 \right) + 2\mu_s^2 \right]}{k_1 k_2 + n_1 n_2} \Bigg\} \\ & - \frac{\lambda'_2 \left[k_1^2 \left(3k_2^2 + n_2^2 \right) + 4k_1 k_2 n_1 n_2 + n_1^2 \left(k_2^2 + 3n_2^2 \right) \right]}{2(k_1 k_2 + n_1 n_2)}, \\ f = {} & (k_1 n_2 - k_2 n_1) \Bigg\{ \frac{2\lambda + \tilde{\lambda}_{12}}{\sqrt{2} k_0} \\ & + \frac{\lambda'_0 k_0^2 + \lambda'_1 (k_1^2 + n_1^2) + \lambda'_2 \left(k_2^2 + n_2^2 \right) + 2\mu_s^2}{\sqrt{2} k_0 (k_1 k_2 + n_1 n_2)} \Bigg\}, \end{aligned} \tag{8}$$

[1] The scalar potential in this limit has been studied in [20].

with the remaining condition $\partial V/\partial k_1 = 0$ automatically satisfied. We conclude this section pointing out the interesting dependence of f on the VEV combination $(k_1 n_2 - k_2 n_1)$. This means that we can interpret the last relation of Eq. (8) as a statement that a small parameter f dynamically induces a small non-zero value for $(k_1 n_2 - k_2 n_1)$, assuming a suitable range of parameters for which the term in curly brackets is non-vanishing.

3. Neutrino masses

After spontaneous symmetry breaking, the Dirac neutrino mass matrix becomes

$$-\mathcal{L}_{\text{mass}} = \frac{1}{\sqrt{2}} \begin{pmatrix} \bar{\nu}_L & \bar{N}_L \end{pmatrix} \begin{pmatrix} y_1 k_1 + \tilde{y}_1 k_2 & y_2 k_1 + \tilde{y}_2 k_2 \\ y_1 n_1 + \tilde{y}_1 n_2 & y_2 n_1 + \tilde{y}_2 n_2 \end{pmatrix} \begin{pmatrix} S_R \\ \tilde{S}_R \end{pmatrix} + \text{h.c.}, \tag{9}$$

where $y_{1,2}$ and $\tilde{y}_{1,2}$ are 3×3 Yukawa matrices. The light neutrino masses can be readily estimated in the one family approximation, in which the neutrino mass matrix is diagonalized by a bi-unitary transformation $\mathcal{M}_{\text{diag}} = U_\nu^\dagger \mathcal{M} U_S$, with

$$U_\alpha \approx \begin{pmatrix} \cos\theta_\alpha & \sin\theta_\alpha \\ -\sin\theta_\alpha & \cos\theta_\alpha \end{pmatrix}, \qquad \alpha = \nu,\ S, \tag{10}$$

and

$$\tan 2\theta_\nu = -\frac{2\left[k_1 n_1 \left(y_1^2 + y_2^2\right) + (k_1 n_2 + k_2 n_1)(\tilde{y}_1 y_1 + \tilde{y}_2 y_2) + k_2 n_2 \left(\tilde{y}_1^2 + \tilde{y}_2^2\right)\right]}{(k_1^2 - n_1^2)\left(y_1^2 + y_2^2\right) + 2(k_1 k_2 - n_1 n_2)(\tilde{y}_1 y_1 + \tilde{y}_2 y_2) + (k_2^2 - n_2^2)\left(\tilde{y}_1^2 + \tilde{y}_2^2\right)},$$
$$\tan 2\theta_S = -\frac{2\left[\left(k_1^2 + n_1^2\right) y_1 y_2 + (k_1 k_2 + n_1 n_2)(\tilde{y}_1 y_2 + \tilde{y}_2 y_1) + (k_2^2 + n_2^2)\tilde{y}_1 \tilde{y}_2\right]}{\left(k_1^2 + n_1^2\right)(y_1^2 - y_2^2) + 2(k_1 k_2 + n_1 n_2)(\tilde{y}_1 y_1 - \tilde{y}_2 y_2) + (k_2^2 + n_2^2)(\tilde{y}_1^2 - \tilde{y}_2^2)}, \tag{11}$$

yielding eigenstates with masses

$$m_\mp = \frac{1}{2}\sqrt{\left(A \mp \sqrt{A^2 - 4B^2}\right)},$$
$$A = \left(k_1^2 + n_1^2\right)\left(y_1^2 + y_2^2\right) + 2(k_1 k_2 + n_1 n_2)(\tilde{y}_1 y_1 + \tilde{y}_2 y_2) + (k_2^2 + n_2^2)\left(\tilde{y}_1^2 + \tilde{y}_2^2\right),$$
$$B = (k_1 n_2 - k_2 n_1)(y_1 \tilde{y}_2 - y_2 \tilde{y}_1). \tag{12}$$

Thus, up to corrections of $\mathcal{O}(B^2/A^{3/2})$, the resulting masses for the light and heavy neutrinos are

$$m_{\text{light}} \approx \frac{|(k_1 n_2 - k_2 n_1)(y_1 \tilde{y}_2 - y_2 \tilde{y}_1)|}{\sqrt{2\left[(k_1^2 + n_1^2)\left(y_1^2 + y_2^2\right) + 2(k_1 k_2 + n_1 n_2)(\tilde{y}_1 y_1 + \tilde{y}_2 y_2) + (k_2^2 + n_2^2)\left(\tilde{y}_1^2 + \tilde{y}_2^2\right)\right]}},$$
$$m_{\text{heavy}} \approx \sqrt{\frac{1}{2}\left[(k_1^2 + n_1^2)\left(y_1^2 + y_2^2\right) + 2(k_1 k_2 + n_1 n_2)(\tilde{y}_1 y_1 + \tilde{y}_2 y_2) + (k_2^2 + n_2^2)\left(\tilde{y}_1^2 + \tilde{y}_2^2\right)\right]}. \tag{13}$$

The mixing angle θ_ν is small by virtue of the VEV hierarchy $k_0, k_1, k_2 \ll n_1, n_2$.[2] For large n_1, n_2 VEVs one clearly obtains the standard seesaw behavior, in which the small neutrino mass emerges from the parameters governing the scale characterizing the mass of the messenger particle, in this case a heavy scalar boson, see Fig. 1.

However Eq. (13) contains further crucial information, namely the fact that the light neutrino mass is determined by the scale $(k_1 n_2 - k_2 n_1)$, the same combination of VEVs found in the last relation in Eq. (8). The smallness of the neutrino mass can then be understood as a consequence of the interplay between the f term in the scalar potential and $V_{\mathcal{L}}$. The presence of a small f parameter (quantifying the amount of $\mathcal{L}$ violation in the scalar sector) and large quartic couplings enforces a nearly parallel dynamical alignment for the ϕ_2 and ϕ_3 VEVs, which in turn can lead to a tiny mass for the active neutrino even without imposing a large hierarchy among the $k_{1,2}$ and $n_{1,2}$ scales.

[2] Notice that θ_S becomes maximal in the limit $y_1 \to y_2$, $\tilde{y}_1 \to \tilde{y}_2$.

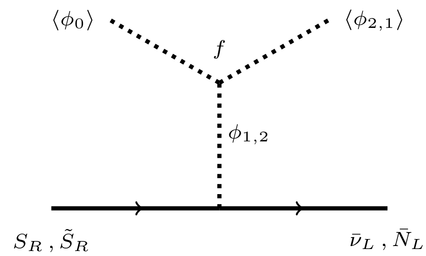

Fig. 1. Type-II-like dynamical seesaw mechanism for Dirac neutrino mass.

Written in terms of f, and using the approximation $n_{1,2} \sim n \gg k_{1,2} \sim k, k_0, |\mu_s|$, the mass of the light neutrino simplifies to

$$m_{\text{light}} \approx \frac{|f k_0 (y_1 \tilde{y}_2 - y_2 \tilde{y}_1)|}{|\tilde{\lambda}_{12} + \lambda_1' + \lambda_2' + 2\lambda| n \sqrt{\left(y_1 + \tilde{y}_1\right)^2 + \left(y_2 + \tilde{y}_2\right)^2}}. \tag{14}$$

Assuming $|\tilde{\lambda}_{12} + \lambda_1' + \lambda_2' + 2\lambda| \sim \mathcal{O}(1)$ in the above equation, we can see that the resulting mass is potentially suppressed by three different sources: (i) the factor k_0/n; (ii) the small scale f associated to the $\mathcal{L}$-symmetry protection, and (iii) the determinant-like Yukawa combination $(y_1 \tilde{y}_2 - y_2 \tilde{y}_1)$. We conclude this section with an illustrative example of how these three sources can act in synergy: setting $f \sim \mathcal{O}(1)$ keV, $k_0 \sim \mathcal{O}(10^2)$ GeV, $n \sim \mathcal{O}(10)$ TeV, and $y_1 = \tilde{y}_2 = y + \delta$, $y_2 = \tilde{y}_1 = y - \delta$ with $\delta \sim \mathcal{O}(10^{-2})$, a naturally small neutrino mass $m_{\text{light}} \sim \mathcal{O}(10^{-1})$ eV is obtained without the need to invoke superheavy physics. The above estimate of light neutrino makes use of the one family approximation. However, generalization to three families is straightforward, using the perturbative block-diagonalization technique developed in [5].

4. Discussion and conclusions

Summarizing, in this letter we have presented a novel mechanism for Dirac neutrino mass generation in the context of $\mathrm{SU(3)_c \otimes SU(3)_L \otimes U(1)_X}$ based models. The light neutrino mass is induced at tree-level *a la seesaw* as indicated in Fig. 1. Its smallness is guaranteed by three independent features:

1. The suppression imposed by the ratio of the $\mathrm{SU(2)_L}$ and $\mathrm{SU(3)_L}$ breaking scales, which by itself can account for the small neutrino mass by invoking a large hierarchy among these scales as in the standard high-scale seesaw mechanism.
2. The dynamical alignment of the VEVs, induced by the presence of a small trilinear term in the scalar potential, which is ultimately related to the smallness of $\mathcal{L}$ violation encoded by the characteristic scale f. The presence of this key ingredient makes it possible to have a low-scale realization of our seesaw mechanism, at an energy scale accessible to the current LHC experiments.
3. The additional suppression provided by the peculiar dependence of the neutrino mass on the Yukawa coefficients of the model, which favors a small neutrino mass in the case of a Yukawa alignment analogous to the one displayed by the VEVs. Potentially this might have a dynamical origin say, in string theories.

Before closing let us note that, as sketched so far, the model provides a low-scale realization of the seesaw mechanism for neutrinos, where the symmetry "protecting" the neutrino mass is $\mathcal{L}$. As such it focuses mainly on the leptons. A realistic pattern of quark masses requires extra scalars beyond those in Table 1. Moreover, alternative variants of the idea are possible, changing the nature of the "protecting" symmetry, this will be considered elsewhere. Typically, this class of models brings in a very rich phenomenology with new quarks, new gauge bosons as well as new scalars, all of them lurking within the range that can be explored in flavor studies and the LHC. Its detailed study lies beyond the scope of this letter. In particular the model can fit the recent hint for a di-photon resonance in a natural way [21,22].

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