Models for the diphoton excess

A large variety of models has been proposed to explain the diphoton excess at 750 GeV. We have selected and implemented several possible models in SARAH as dscribed in Ref. [1]. Our selection is not exhaustive, but we have tried to implement a sufficient cross-section which are representative of many of the ideas put forward in the context of renormalisable models. These are the ones that SARAH can handle. Their description is organised in the subsections that follow. The main purpose of this document is to give the essential information about the implementation, so that it can be used and modified by interested users. We do not intend to present a thorough study of the models, however in some cases we do make some observations about the motivation and validity of the models.

All model files are available for download at

http://sarah.hepforge.org/Diphoton_Models.tar.gz

and an overview of all implemented models is given in Tabs. 1 and 2. In case of questions, comments or bug reports concerning these models, please, send an e-mail to diphoton-tools@cern.ch which includes all authors. If you use one of the model files presented here, please cite the according SARAH references along with Ref. [1].

1 Toy models

The simplest ideas proposed to explain the diphoton excess extend the SM by a scalar singlet and vector-like fermions, which serve the purpose of enhancing the diphoton rate, and – in the case of coloured states – also the production via gluon fusion. An enhancement of gluon fusion seems to be necessary because a production of the resonance purely by photon fusion is in some tension with 8 TeV data: the increase in the cross section from 8 to 13 TeV is just a factor 2, while a factor of 5 would be needed to make the results from LHC run-I and II compatible. To explore the multitude of possibilities we first consider toy models: they do not contain all possible couplings of the vector-like fermions to SM matter, but are rather engineered to allow one to easily explore the effects of different representations of vector-like matter on the diphoton rate and on the relevant partial widths. Then, from section 1.1 on, we consider complete models, containing all the operators consistent with both field content and symmetries.

Toy models with vector-like fermions

- **Reference:** [2, 43–45]

Model	Name	Section	Ref.	
Toy models with vector-like fermions				
CP-even singlet	`SM+VL/CPevenS`	1		
CP-odd singlet	`SM+VL/CPoddS`	1		
Complex singlet	`SM+VL/complexS`	1		
Models based on the SM gauge-group				
Portal dark matter	`SM+VL/PortalDM`	1.2.1	[2, 3]	
Scalar octet	`SM-S-Octet`	1.2.2	[4, 5]	⚠ (1)
$SU(2)$ triplet quark model	`SM+VL/TripletQuarks`	1.2.3	[6]	
Single scalar leptoquark	`LQ/ScalarLeptoquarks`	1.2.4	[7]	
Two scalar leptoquarks	`LQ/TwoScalarLeptoquarks`	1.2.5	[8]	⚠ (3)
Georgi-Machacek model	`Georgi-Machacek`	1.2.6	[9, 10]	
THDM w. colour triplet	`THDM+VL/min-3`	1.3.1	[11]	
THDM w. colour octet	`THDM+VL/min-8`	1.3.1	[11]	
THDM-I w. exotic fermions	`THDM+VL/Type-I-VL`	1.3.2	[12, 13]	
THDM-II w. exotic fermions	`THDM+VL/Type-II-VL`	1.3.2	[12, 13]	
THDM-I w. SM-like fermions	`THDM+VL/Type-I-SM-like-VL`	1.3.3	[14]	
THDM-II w. SM-like fermions	`THDM+VL/Type-II-SM-like-VL`	1.3.3	[14]	
THDM w. scalar septuplet	`THDM/ScalarSeptuplet`	1.3.4	[15, 16]	

Table 1. Part I of the overview of proposed models to explain the diphoton excess which are now available in `SARAH`. The warning (⚠) shows that we found serious problems with the model during the implementation. The reasons are as follows. (1): the model is in conflict with limits from $S \to jj$; (2): we changed the quantum numbers and/or the potential because the original model had charge violating interactions; (3): we find disagreement with the diphoton rate as calculated in the original reference. For simplicity, we used the abbreviations `LQ` for `LeptoQuarks` and `U1Ex` for `U1Extensions` .

- **Model names:**
  `SM+VL/CPevenS`
  `SM+VL/CPoddS`
  `SM+VL/complexS`

To begin with we categorise the toy models according to the type of the involved scalar singlet. There are three possibilities: (i) the singlet is a real CP-even scalar, (ii) real CP-odd, or (iii) a complex scalar. Each case is considered in separate `SARAH` model files, where we introduce all possible representations of vector-like fermions. These possibilities, following Tables 3 and 4 of Ref. [43], are shown below in Table 3. This

Model	Name	Section	Ref.	
	$U(1)$ **Extensions**			
Dark $U(1)'$	`U1Ex/darkU1`	2.1.1	[17]	
Hidden $U(1)$	`U1Ex/hiddenU1`	2.1.2	[18]	
Simple $U(1)$	`U1Ex/simpleU1`	2.1.3	[19]	
Scotogenic $U(1)$	`U1Ex/scotoU1`	2.1.4	[20]	⚠(2)
Unconventional $U(1)_{B-L}$	`U1Ex/BL-VL`	2.2.1	[21]	
Sample of $U(1)'$	`U1Ex/VLsample`	2.2.2	[22]	
flavour-nonuniversal charges	`U1Ex/nonUniversalU1`	2.2.3	[23]	
Leptophobic $U(1)$	`U1Ex/U1Leptophobic`	2.2.4	[24]	⚠(1)
Z' mimicking a scalar resonance	`U1Ex/trickingLY`	2.2.5	[25]	
	Non-abelian gauge-group extensions of the SM			
LR without bidoublets	`LRmodels/LR-VL`	3.1.1	[26–28]	⚠(2)
LR with $U(1)_L \times U(1)_R$	`LRmodels/LRLR`	3.1.2	[29]	⚠(2)
LR with triplets	`LRmodels/tripletLR`	3.1.3	[30]	
Dark LR	`LRmodels/darkLR`	3.1.4	[31]	
331 model without exotic charges	`331/v1`	3.2.1	[32]	
331 model with exotic charges	`331/v2`	3.2.2	[33]	
Gauged THDM	`GTHDM`	3.3	[34]	
	Supersymmetric models			
NMSSM with vectorlike top	`NMSSM+VL/VLtop`	4.2	[35]	⚠(1)
NMSSM with **5**'s	`NMSSM+VL/5plets`	4.2.1	[36–38]	
NMSSM with **10**'s	`NMSSM+VL/10plets`	4.3	[36–38]	
NMSSM with **5**'s & **10**'s	`NMSSM+VL/10plets`	4.4	[38]	
NMSSM with **5**'s and RpV	`NMSSM+VL/5plets+RpV`	4.5	[36]	
Broken MRSSM	`brokenMRSSM`	4.6	[39]	
$U(1)'$-extended MSSM	`MSSM+U1prime-VL`	4.7	[40, 41]	
E_6 with extra $U(1)$	`E6MSSMalt`	4.8	[42]	

Table 2. Part II of the overview of proposed models to explain the diphoton excess which are now available in `SARAH`. The warning (⚠) shows that we found serious problems with the model during the implementation. The reasons are as follows. (1): non-perturbative couplings needed to explain diphoton excess; (2): we changed the quantum numbers and/or the potential because the original model had charge violating interactions; (3): we find disagreement with the diphoton rate as calculated in the original reference.

allows one to study combinations of fermion representations or individual choices by giving unwanted fermion representations a mass high enough to effectively decouple them from the model [1]. All mixings between the extra fermions and SM fermions are neglected through the assumption of a discrete $\mathbb{Z}_2$ symmetry. Of course in a realistic model the mixings have to be taken into account, as they allow the necessary decays of the coloured vector-like fermions into SM particles.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$\mathbb{Z}_2$	Ref.
S	1	1	1	0	$+$	
Ψ_{F_1}	1	$\mathbf{3}$	$\mathbf{2}$	$\frac{7}{6}$	$-$	
Ψ_{F_2}	1	$\mathbf{3}$	$\mathbf{3}$	$\frac{2}{3}$	$-$	
Ψ_{F_3}	1	$\mathbf{3}$	$\mathbf{2}$	$-\frac{5}{6}$	$-$	
Ψ_{F_4}	1	$\mathbf{3}$	$\mathbf{3}$	$-\frac{1}{3}$	$-$	
Ψ_{F_5}	1	$\mathbf{3}$	1	$\frac{2}{3}$	$-$	[44, 45]
Ψ_{F_6}	1	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$	$-$	
Ψ_{F_7}	1	$\mathbf{3}$	1	$-\frac{1}{3}$	$-$	[45]
Ψ_{F_8}	1	1	1	1	$-$	
Ψ_{F_9}	1	1	$\mathbf{2}$	$-\frac{3}{2}$	$-$	
$\Psi_{F_{10}}$	1	1	$\mathbf{3}$	1	$-$	
$\Psi_{F_{11}}$	1	1	$\mathbf{2}$	$-\frac{1}{2}$	$-$	
$\Psi_{F_{12}}$	1	1	$\mathbf{3}$	0	$-$	
$\Psi_{F_{13}}$	1	$\mathbf{3}$	1	$\frac{5}{3}$	$-$	[2]

Table 3. Extra particle content of the toy models. S is either the CP-even, CP-odd or complex scalar. The various fermions $\Psi_{F_i} \equiv \Psi_{F_{iL}}$ each come with a right-handed partner $\Psi_{F_{iR}}$ with opposite quantum numbers. These models are based on the collection given in Ref. [43], while the last column contains other works where fermions in these specific representations are used. All SM particles have charge '+' under the additional $\mathbb{Z}_2$ symmetry.

[1] This option has to be used carefully when including loop corrections to the mass spectrum.

We write the scalar potentials for the three different types of scalars as

$$\begin{aligned}
V_{\text{even}} &= \frac{1}{2}m_S^2 S^2 + \frac{1}{4}\lambda_S S^4 - \mu^2|H|^2 + \frac{1}{2}\lambda_H|H|^4 + \frac{1}{2}\lambda_{HS}S^2|H|^2 \\
&\quad + \kappa_{HS}S|H|^2 + \frac{1}{3}\kappa_S S^3\,, \qquad (1.1a)\\
V_{\text{odd}} &= \frac{1}{2}m_S^2|S|^2 + \frac{1}{4}\lambda_S|S|^4 - \mu^2|H|^2 + \frac{1}{2}\lambda_H|H|^4 + \frac{1}{2}\lambda_{HS}|S|^2|H|^2 \\
&\quad + \left(i\kappa_{HS}S|H|^2 + i\frac{1}{3}\kappa_S S|S|^2 + \text{h.c.}\right)\,, \qquad (1.1b)\\
V_{\text{complex}} &= m_S^2|S|^2 + \frac{1}{2}\lambda_S|S|^4 - \mu^2|H|^2 + \frac{1}{2}\lambda_H|H|^4 + \lambda_{HS}|S|^2|H|^2 \\
&\quad + \left(\kappa_{HS}S|H|^2 + \frac{1}{3}\kappa_S S|S|^2 + \text{h.c.}\right)\,. \qquad (1.1c)
\end{aligned}$$

The Yukawa interactions are given by

$$-\mathcal{L}_Y = \mathcal{L}_Y^{\text{SM}} + \sum_j \left(m_{F_j}\overline{\Psi_{F_jL}}\Psi_{F_jR} + Y_{F_j}S\,\overline{\Psi_{F_jL}}\Psi_{F_jR}\right) + \text{h.c.}\,. \qquad (1.2)$$

In the Lagrangian above one should substitute the expression for the relevant scalar field

$$\begin{aligned}
S_{\text{even}} &= v_S + \phi_S\,, \qquad \text{where} \quad \langle S\rangle = v_S\,, \qquad (1.3a)\\
S_{\text{odd}} &= i\sigma_S\,, \qquad (1.3b)\\
S_{\text{complex}} &= \frac{1}{\sqrt{2}}\left(v_S + \phi_S + i\sigma_S\right)\,. \qquad (1.3c)
\end{aligned}$$

Note that imposing CP conservation forces κ_{HS} and κ_S to vanish in the CP-odd potential. For both the CP-even and complex singlet models the CP-even component ϕ_S mixes with the neutral Higgs field ϕ_h at tree-level if $\kappa_{HS} \neq 0$. As discussed in [1], even if one sets $\kappa_{HS} = 0$ mixing between the CP-even states is induced at the loop level.

1.1 Models based on the SM gauge-group

We now turn our attention to complete models that have been proposed to explain the 750 GeV diphoton excess. To begin with we consider models that are based on the SM gauge group, with or without additional global symmetries. We divide the possible models into two main categories: (i) models with a SM-like Higgs sector and (ii) Two-Higgs-doublet type models.

1.2 Singlet extensions with vector-like fermions

1.2.1 Portal dark matter model

- **Reference:** [2, 3]
- **Model name:** `SM+VL/PortalDM`

This model proposes that the resonance is produced by a 750 GeV real scalar singlet S, with the diphoton rate boosted through the introduction of vector-like quarks coupling to the new scalar singlet. In this model we have three possible options for the representation of the new vector-like matter. These choices are: (i) the addition of a vector-like up-type quark pair $t'_{L/R}$ [3], (ii) in addition to the vector-like up-type quark pair, a vector-like quark doublet pair $Q'_{L/R}$ is introduced [3] and finally, (iii) the addition of only the vector-like pair $X_{L/R}$, also triplet under $SU(3)_C$ but with exotic hypercharge [2].

The model also introduces a new real scalar singlet S_{DM} and an additional fermionic singlet Ψ_{DM} as DM candidates, with a $\mathbb{Z}_2$ symmetry to stabilise them. The particle content beyond the SM fields is given in Table 4. In order to avoid mixing with the SM quarks, the fields $t'_{L/R}$ and $Q'_{L/R}$ are also odd under the $\mathbb{Z}_2$.

The user can choose between the three model types by setting the couplings of the unwanted fields to zero and choosing their masses to be very large (for example, 10^{12} GeV) to decouple them. Originally, the fermionic dark matter is absent from the models described in [3]. The exact settings are given below.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$\mathbb{Z}_2$
S	1	$\mathbf{1}$	$\mathbf{1}$	0	$+$
S_{DM}	1	$\mathbf{1}$	$\mathbf{1}$	0	$-$
Ψ_{DM}	1	$\mathbf{1}$	$\mathbf{1}$	0	$-$
X	1	$\mathbf{3}$	$\mathbf{1}$	$\frac{5}{3}$	$-$
t'	1	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	$-$
Q'	1	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$	$-$

Table 4. Extra particle content of the portal DM model. The top/bottom part of the table corresponds to the new scalar/fermionic degrees of freedom. All additional fermionic degrees of freedom are vector-like fermions.

The scalar potential for these models reads

$$V = -\mu^2|H|^2 + \frac{1}{2}\lambda_H\,|H|^4 + \frac{1}{2}M_S^2\,S^2 + \frac{1}{3}\kappa_S\,S^3 + \frac{1}{4}\lambda_S\,S^4 + \frac{1}{2}M_{S_{\rm DM}}^2\,S_{\rm DM}^2 + \frac{1}{4}\lambda_{S_{\rm DM}}\,S_{\rm DM}^4 \\ + \kappa_{HS}\,|H|^2 S + \lambda_{HS}\,|H|^2 S^2 + \lambda_{HS_{\rm DM}}\,|H|^2 S_{\rm DM}^2 + \kappa_{SS_{\rm DM}} S S_{\rm DM}^2 + \lambda_{SS_{\rm DM}} S^2 S_{\rm DM}^2\,, \quad (1.4)$$

whereas the three model variants lead to three distinct forms for the Yukawa interactions, given by:

$$-\mathcal{L}_Y^{\rm I} = \mathcal{L}_Y^{\rm SM} + (m_{t'} + Y_{St'}S)\,\overline{t'}_L\,t'_R + Y_{S_{\rm DM}t'}S_{\rm DM}\overline{t'_L}u_R + \text{h.c.}\,, \quad (1.5a)$$

$$-\mathcal{L}_Y^{\rm II} = \mathcal{L}_Y^{\rm I} + (m_{Q'} + Y_{SQ'}S)\,\overline{Q'}_L Q'_R \\ + Y_{Q'_1}\overline{Q'}_L H t'_R + Y_{Q'_2}\overline{Q'}_R \widetilde{H} t'_L + Y_{S_{\rm DM}Q'}S_{\rm DM}\overline{Q_L}Q'_R + \text{h.c.}\,, \quad (1.5b)$$

$$-\mathcal{L}_Y^{\rm III} = \mathcal{L}_Y^{\rm SM} + (m_{\rm DM} + \kappa S)\,\overline{\Psi_{\rm DM}}_L\Psi_{{\rm DM}R} + (m_X + Y_X\,S)\,\overline{X}_L X_R + \text{h.c.}\,, \quad (1.5c)$$

where $\widetilde{H} = i\sigma_2 H^*$. In the model variants I and II, the vectorlike quarks decay into SM quarks and the scalar dark matter candidate $S_{\rm DM}$ via the couplings $Y_{S_{\rm DM}Q'}$ and $Y_{S_{\rm DM}t'}$. In model III, in turn, X is stable at the level of the Lagrangian and could only decay through higher-dimensional operators which are not included here.

The symmetry breaking pattern of the models is that of the SM, where the neutral component of the Higgs field acquires a VEV, plus the VEV of the scalar singlet S

$$S = v_S + \phi_S\,, \qquad \text{where} \quad \langle S\rangle = v_S\,. \quad (1.6)$$

In general, ϕ_S mixes with the SM Higgs.

As mentioned previously, the user can choose between the three different models through the following parameter choices:

- **Model I:** $Y_{SQ'} = Y_{Q'_i} = Y_{S_{\rm DM}Q'} = Y_X = 0$ and $m_{Q'} = m_X = 10^{12}\,\text{GeV}$
- **Model II:** $Y_X = 0$ and $m_X = 10^{12}\,\text{GeV}$
- **Model III:** $Y_{SQ'} = Y_{Q'_i} = Y_{S_{\rm DM}Q'} = Y_{St'} = Y_{S_{\rm DM}} = 0$ and $m_{Q'} = m_{t'} = 10^{12}\,\text{GeV}$

1.2.2 Scalar octet extension

- **Reference:** [4, 5]
- **Model name:** `SM-S-Octet`

A charged scalar colour octet O coupled to a scalar singlet S was proposed in Refs. [4, 5]. Here the singlet is the 750 GeV candidate, while the octet enters the loops that contribute to the generation of the couplings of the singlet to the gauge bosons. While

Ref. [5] considers a toy model involving only the term $S\,|O|^2$, Ref. [4] takes the singlet extended Manohar-Wise model [46]. For the `SARAH` implementation we have used the full model. However, since the cubic and quartic terms in O do not play a significant role, they are turned off by default in the `SARAH` model file.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$
S	1	$\mathbf{1}$	$\mathbf{1}$	0
O	1	$\mathbf{8}$	$\mathbf{2}$	$\frac{1}{2}$

Table 5. Extra scalar field content of the octet extended SM.

The extra particle content with respect to the SM is a real singlet S and a scalar color octet O which is also charged under $SU(2)_L \times U(1)_Y$, see Table 5. The isospin components of O are

$$O^A = \begin{pmatrix} O^{+\,A} \\ O^{0\,A} \end{pmatrix}, \tag{1.7}$$

where $A = 1, \ldots, 8$ is the adjoint colour index. The full scalar potential reads

$$\begin{aligned} V = &\frac{1}{2} m_S^2 S^2 + \lambda_S S^4 - \mu^2 |H|^2 + \lambda_H |H|^4 + \kappa_1 S^2 |H|^2 + 2 m_O^2 \mathrm{Tr}(O^\dagger O) + \kappa_2 S^2 \mathrm{Tr}(O^\dagger O) \\ &+ \lambda_1 |H|^2 \mathrm{Tr}(O^\dagger O) + \lambda_2 H_i^\dagger H_j \mathrm{Tr}(O_j^\dagger O_i) + \lambda_6 \mathrm{Tr}(O^\dagger O O^\dagger O) + \lambda_7 \mathrm{Tr}(O_i^\dagger O_j O_j^\dagger O_i) \\ &+ \lambda_8 \mathrm{Tr}(O^\dagger O)^2 + \lambda_9 \mathrm{Tr}(O_i^\dagger O_j) \mathrm{Tr}(O_j^\dagger O_i) + \lambda_{10} \mathrm{Tr}(O_i O_j) \mathrm{Tr}(O_i^\dagger O_j^\dagger) + \lambda_{11} (O_i O_j O_j^\dagger O_i^\dagger) \\ &+ \Big(\lambda_3 H_i^\dagger H_j^\dagger \mathrm{Tr}(O_i O_j) + \lambda_4 H_i^\dagger \mathrm{Tr}(O_j^\dagger O_j O_i) + \lambda_5 H_i^\dagger \mathrm{Tr}(O_j^\dagger O_i O_j) + \text{h.c.} \Big)\,. \end{aligned} \tag{1.8}$$

Electroweak symmetry-breaking (EWSB) is driven by the VEV of the neutral component of the SM Higgs doublet, which can be decomposed as

$$H^0 = \frac{1}{\sqrt{2}} \left(v + \phi_H + i\, \sigma_H \right)\,. \tag{1.9}$$

Here $\phi_H \equiv h$ is the Higgs boson, to be identified with the 125 GeV state discovered at the LHC. Similarly, the singlet S receives a VEV, and the neutral component of the octet is split into its CP-even and CP-odd eigenstates:

$$S = v_S + \phi_S\,, \qquad O^0 \to \frac{1}{\sqrt{2}} \left(O^R + i\, O^I \right)\,. \tag{1.10}$$

We will now briefly discuss the parameter space of the model in order to justify our choice of input parameters. First, we consider the tadpole equations, which can be automatically derived by `SARAH`. Their solution for μ^2 and κ_1 is

$$\begin{aligned}\mu^2 =& -\frac{1}{v^2}(\lambda_H v^4 - m_S^2 v_S^2 - 4\lambda_S v_S^4)\,,\\ \kappa_1 =& -\frac{1}{v^2}(m_S^2 + 4\lambda_S v_S^2)\,.\end{aligned} \tag{1.11}$$

The tree-level mass matrix for the CP-even neutral scalars in the (ϕ_H, ϕ_S) basis is given by

$$\begin{aligned}\mathcal{M}^2 =& \begin{pmatrix} \mu^2 + 3\lambda_H v^2 + \kappa_1 v_S^2 & 2\kappa_1 v v_S \\ 2\kappa_1 v v_S & m_S^2 + \kappa_1 v^2 + 12\lambda_S v_S^2 \end{pmatrix}\\ =& \begin{pmatrix} 2\lambda_H v^2 & -\frac{2v_S}{v}(m_S^2 + 4\lambda_S v_S^2) \\ -\frac{2v_S}{v}(m_S^2 + 4\lambda_S v_S^2) & 8\lambda_S v_S^2 \end{pmatrix}.\end{aligned} \tag{1.12}$$

We note that, in general, there is singlet-doublet mixing. There are two reasons to consider a small singlet-doublet mixing angle, θ. First, the stringent constraints derived from Higgs physics measurements, and second, the required suppressed decay widths into Higgses, W's and Z's in order to fit the diphoton signal – indeed in [4] values of $\sim 10^{-2}$ were found to be required. If we have a small mixing angle, then we can write

$$\mathcal{M}^2 \sim \begin{pmatrix} m_h^2 & s_\theta c_\theta (m_h^2 - m_{750}^2) \\ s_\theta c_\theta (m_h^2 - m_{750}^2) & m_{750}^2 \end{pmatrix}. \tag{1.13}$$

This implies $\lambda_S > 0$, but also

$$\mu^2 \simeq -\frac{1}{2} m_h^2 + \frac{v_S^2}{v^2}(m_S^2 + \frac{1}{2} m_{750}^2)\,. \tag{1.14}$$

However, we also have $v_S^2 \sim m_{750}^2 / 8\lambda_S$, and so

$$\mu^2 \simeq -\frac{1}{2} m_h^2 + \frac{1.2}{\lambda_S}(m_S^2 + \frac{1}{2} m_{750}^2)\,. \tag{1.15}$$

We thus require a tachyonic m_S^2 for the SM Higgs mass condition:

$$m_S^2 \simeq -\frac{1}{2} m_{750}^2 + \frac{\lambda_S}{1.2}(\mu^2 + \frac{1}{2} m_h^2) \lesssim -(500 \text{ GeV})^2 \tag{1.16}$$

where in the last step we have taken $\lambda_S = 1.2$, a rather large value. If we want $\kappa_1 \sim -1$ then we require $m_S^2 \sim -(600 \text{ GeV})^2$. On the other hand, from the second tadpole equation we have

$$m_S^2 = -\kappa_1 v^2 - \frac{1}{2} m_{750}^2, \tag{1.17}$$

which, if we require $|\kappa_1| < 2$, gives

$$-(630 \text{ GeV})^2 \leq m_S^2 \leq -(400 \text{ GeV})^2\,, \tag{1.18}$$

so putting these together we find the narrow window

$$-(630 \text{ GeV})^2 \leq m_S^2 \leq -(500 \text{ GeV})^2\,. \tag{1.19}$$

Alternative implementation in SARAH

The above discussion suggests to use a different choice for the input parameters of the model in our SARAH implementation: ideally we would like the particle masses, the mixing and only dimensionless couplings to be the inputs. We shall take the input parameters to be

$$m_h, m_{750}, s_\theta, \lambda_S. \tag{1.20}$$

In terms of these the other parameters are determined to be

$$\begin{aligned}
\lambda_H =& \frac{c_\theta^2 m_h^2 + s_\theta^2 m_{750}^2}{2v^2}, \qquad v_S^2 = \frac{c_\theta^2 m_{750}^2 + s_\theta^2 m_h^2}{8\lambda_S} \\
m_S^2 =& -\kappa_1 v^2 - \frac{1}{2} m_{750}^2, \qquad \kappa_1 v v_S = s_\theta c_\theta (m_h^2 - m_{750}^2) \\
\to \kappa_1 =& \frac{\sqrt{2\lambda_S} s_\theta c_\theta (m_h^2 - m_{750}^2)}{v\sqrt{(c_\theta^2 m_{750}^2 + s_\theta^2 m_h^2)}} \simeq -4.3 \times s_\theta \sqrt{\lambda_S}
\end{aligned} \tag{1.21}$$

The exact version of these equations is implemented in SARAH and can be selected using the `InputFile`$\to$`"SPheno_diphoton.m"` option in `MakeAll` or `MakeSPheno`.

Octet masses

One further input is taken in [4]: the physical mass of the octet scalars. These are given in terms of the Lagrangian parameters as:

$$\begin{aligned}
m_{O_r^0}^2 =& m_O^2 + \kappa_2 v_S^2 + \frac{v^2}{2}(\lambda_1 + \lambda_2 + 2\,\mathrm{Re}(\lambda_3)) \\
m_{O_i^0}^2 =& m_O^2 + \kappa_2 v_S^2 + \frac{v^2}{2}(\lambda_1 + \lambda_2 - 2\,\mathrm{Re}(\lambda_3)) \\
m_{O^+}^2 =& m_O^2 + \kappa_2 v_S^2 + \frac{1}{2}\lambda_1 v^2
\end{aligned} \tag{1.22}$$

The values of λ_i are taken to be small and equal in order for the octets to have similar masses, but since this is not the general case, we do not impose this choice in SARAH.

The choice in that paper does however hide the possibility of tachyonic m_O^2 (and hence possible charge/colour breaking minima) – indeed, if we insist that $m_O^2 > 0$ we have a lower bound on the masses of

$$m^2_{O^{0,+}} > \frac{\kappa_2}{8\lambda_S} m^2_{750}. \tag{1.23}$$

Clearly this is violated for $m_{O^{0,+}} = 600$ GeV when $\kappa_2 \sim 1, \lambda_S \ll 1$. On the other hand, this does not guarantee a problem.

The desired vacuum has energy

$$\begin{aligned} V_0 =& \frac{m_S^2 v_S^2}{2} - \frac{\lambda_H}{4} v^4 + \lambda_S v_S^4 \\ \simeq& -\frac{1}{8} v^2 m_h^2 - v_S^2 (\frac{1}{2}\kappa_1 v^2 + \frac{1}{4} m^2_{750} - \frac{m^2_{750}}{8}) \\ \simeq& -\frac{1}{8} v^2 m_h^2 - \frac{m^2_{750}}{8\lambda_S}(-2 s_\theta \sqrt{\lambda_S} + \frac{m^2_{750}}{8}) \end{aligned} \tag{1.24}$$

If we instead concentrate on the potential terms containing the octets, where only one component develops a VEV, we find

$$\begin{aligned} V(O^R) =& \frac{1}{2}(O^R)^2 \left[m_O^2 + \frac{1}{8}(\lambda_9 + \lambda_{10} + \frac{1}{9}\lambda_6 + \frac{1}{9}\lambda_7 + \frac{1}{9}\lambda_{11})(O^R)^2 \right] \\ V(O^I) =& \frac{1}{2}(O^I)^2 \left[m_O^2 + \frac{1}{8}(\lambda_9 + \lambda_{10} + \frac{1}{9}\lambda_6 + \frac{1}{9}\lambda_7 + \frac{1}{9}\lambda_{11})(O^I)^2 \right] \\ V(O^+) =& |O^+|^2 \left[m_O^2 + \frac{1}{4}(\lambda_9 + \lambda_{10} + \frac{1}{9}\lambda_6 + \frac{1}{9}\lambda_7 + \frac{1}{9}\lambda_{11})|O^+|^2 \right] \end{aligned} \tag{1.25}$$

Arranging for the additional minimum of the potential to be higher than the colour-breaking one then places a *lower* bound on the octet self-couplings, but for the phenomenology of the diphoton excess – when we neglect loop corrections to the mass of the octet – they play no other role.

Comments on fitting the excess

In [4] the authors find that the diphoton excess can be easily fit with octets at 600 or 1000 GeV and $\kappa_2 \sim 1.5$ or 4.5, respectively. The scenario involves merely the simplifying assumption $\lambda_1 = \lambda_2 = \lambda_3$ so that the octets are of approximately equal mass. The ratio between the digluon and diphoton decay rates is then

$$\frac{\Gamma(S \to gg)}{\Gamma(S \to \gamma\gamma)} \simeq \frac{9}{2}\frac{\alpha_s^2}{\alpha^2}. \tag{1.26}$$

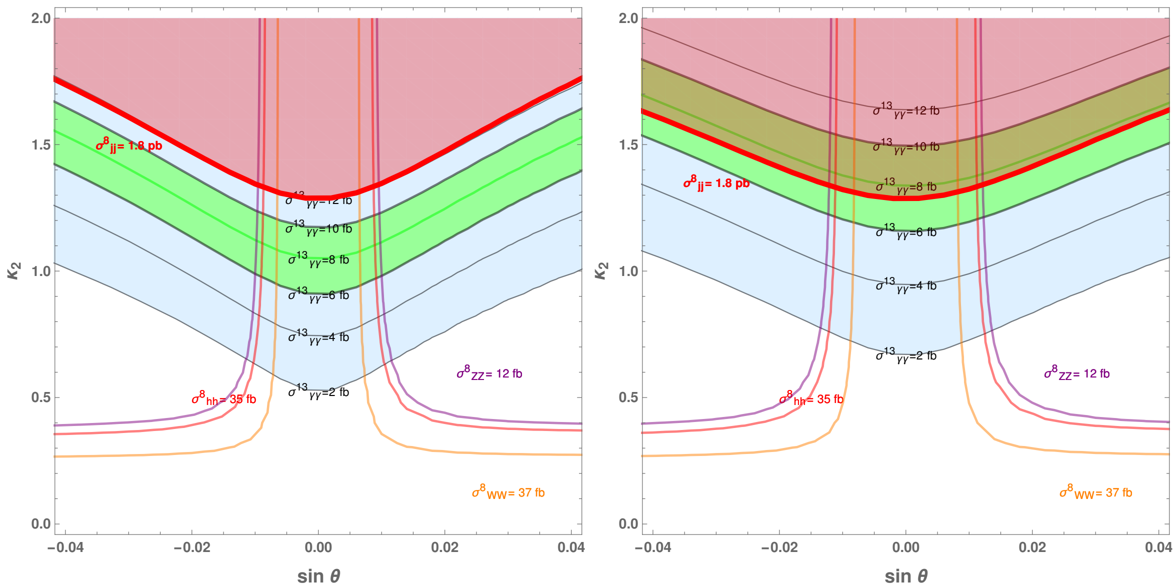

Figure 1. Scan over sine of Higgs-singlet mixing angle θ and κ_2 for octet masses of 600 GeV, $\lambda_S = 0.07$ corresponding to $v_S \simeq 1000$ GeV. The contours show the 750 GeV resonance production cross-sections σ^X_{YY} at energy X TeV decaying into channel YY. On the left plot, only leading order contributions to the decays are used; on the right, all corrections up to N^{3}LO available in `SARAH` are included.

In [4] this is quoted as $\simeq 715$. In `SARAH`, before any NLO corrections are applied, the running of the Standard Model gauge couplings yields $\alpha_s(750 \text{ GeV}) = 0.091$ and we use $\alpha(0) \simeq 137^{-1}$, giving a ratio of 700, in good agreement. However, when we include corrections up to N^{3}LO, this ratio rises to 1150, putting the model near the boundary of exclusion due to dijet production at 8 TeV. These differences are illustrated in plots produced from `SARAH`/`SPheno` in figure 1. To produce these plots, all branching ratios/widths are calculated in `SPheno`, as is the production cross-section of the resonance at 8 TeV. To calculate 13 TeV cross-sections the 8 TeV cross-section was rescaled by the parton luminosity factor for gluons of 4.693.

1.2.3 Vector-like $SU(2)$ triplet quark model

- **Reference:** [6]
- **Model name:** `SM+VL/TripletQuarks`

The model introduced in [6] considers a singlet scalar S as the candidate for the 750 GeV resonance. In order to produce the singlet via gluon fusion at the LHC, the model is further extended with the introduction of vector-like quarks, triplet under $SU(2)_L$. Moreover, the singlet is charged under a $\mathbb{Z}_2$ parity, although this is softly broken by the vector-like quark mass terms.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$\mathbb{Z}_2$
S	1	$\mathbf{1}$	$\mathbf{1}$	0	$-$
F_{1L}	1	$\mathbf{3}$	$\mathbf{3}$	$\frac{2}{3}$	$-$
F_{1R}	1	$\mathbf{3}$	$\mathbf{3}$	$\frac{2}{3}$	$+$
F_{2L}	1	$\mathbf{3}$	$\mathbf{3}$	$-\frac{1}{3}$	$-$
F_{2R}	1	$\mathbf{3}$	$\mathbf{3}$	$-\frac{1}{3}$	$+$

Table 6. Extra scalar/fermionic degrees of freedom shown in the top/bottom. All SM particles are even under the imposed discrete symmetry.

This model is based on the SM gauge symmetry, extended with a $\mathbb{Z}_2$ parity. The fermionic and scalar particle content is summarized in Table 6. The vector-like $SU(2)_L$ triplet quarks can be expressed in 2×2 matrix notation as

$$F_1 = \begin{pmatrix} U_1/\sqrt{2} & X \\ D_1 & -U_1/\sqrt{2} \end{pmatrix}, \qquad F_2 = \begin{pmatrix} D_2/\sqrt{2} & U_2 \\ Y & -D_2/\sqrt{2} \end{pmatrix}. \tag{1.27}$$

Here we see that the $U_{1,2}$ and $D_{1,2}$ components have the same electric charges as the SM up- and down-type quarks, respectively. The components X and Y are exotic coloured states with electric charges $5/3$ and $-4/3$, respectively, and thus they do not mix.

The Yukawa Lagrangian of the model can be written as

$$\begin{aligned} -\mathcal{L}_Y = {} & \mathcal{L}_Y^{\rm SM} + Y_1\, \overline{Q_L}\, F_{1\,R}\, \widetilde{H} + Y_2\, \overline{Q_L}\, F_{2\,R}\, H \\ & + (m_{F_1} + Y_{F_1 S} S)\, \overline{F_1}\, F_1 + (m_{F_2} + Y_{F_2 S} S)\, \overline{F_2}\, F_2 \; + \text{h.c.}\,, \end{aligned} \tag{1.28}$$

whereas the scalar potential is given by

$$V = -\mu^2\, |H|^2 + \lambda_H\, |H|^4 + m_S^2\, S^2 + \frac{1}{4}\lambda_S\, S^4 + \frac{1}{3}\lambda_{HS}\, |H|^2 S^2\,. \tag{1.29}$$

We note that the vector-like mass terms for F_1 and F_2 softly break the $\mathbb{Z}_2$ parity. We assume that the SM Higgs doublet obtains a VEV while the introduced singlet does not, hence

$$H = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + \phi_H + i\,\sigma_H \end{pmatrix}, \qquad S = \phi_S\,. \tag{1.30}$$

The $\mathbb{Z}_2$ discrete symmetry, together with $\langle S \rangle = 0$ ensured by the symmetry, implies that no mixing between the SM Higgs boson ϕ_H and the singlet S appears at tree-level.

1.2.4 Single scalar leptoquark model

- **Reference:** [7]
- **Model name:** `Leptoquarks/ScalarLeptoquarks`

The model introduced in Ref. [7] considers a singlet scalar, S, as candidate for the 750 GeV resonance. It is based on the scalar leptoquark model in Ref. [47], with the addition of a vector-like fermion multiplet χ transforming in an a priori unspecified representation of $SU(2)_L$. The simplest model, which the authors find to work, requires χ to be a $SU(2)_L$ triplet. The model is based on the gauge symmetries of the SM, augmented with either a discrete or gauge symmetry necessary for the DM sector. For the sake of simplicity we choose to realise the model using only a discrete $\mathbb{Z}_2$ symmetry. The particle content beyond the SM is shown in Table 7.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$\mathbb{Z}_2$
S	1	**1**	**1**	0	$+$
ϕ	1	**3**	**1**	$-\frac{1}{3}$	$+$
χ	N_χ	**1**	**3**	Y_χ	$-$

Table 7. Additional scalar and fermionic particles beyond the SM shown in the top and bottom respectively. Note that all SM fields are neutral under the $\mathbb{Z}_2$ symmetry.

The new degrees of freedom are: the scalar leptoquarks ϕ, the gauge singlet scalar S, which is assumed to be real, and finally the vector-like triplet fermions χ. The vector-like $SU(2)_L$ triplet can be expressed in 2×2 matrix notation as

$$\chi = \begin{pmatrix} \frac{\chi^0}{\sqrt{2}} & \chi^+ \\ \chi^- & -\frac{\chi^0}{\sqrt{2}} \end{pmatrix} . \tag{1.31}$$

Due to the discrete $\mathbb{Z}_2$ symmetry the neutral component χ^0 is stable and a suitable DM candidate. In Ref. [7] the hypercharge Y_χ and the number of generations N_χ are left as free parameters. In order to boost the ratio of decay widths $\Gamma_{S\to\gamma\gamma}/\Gamma_{S\to gg}$ they claim that the minimal choice is $N_\chi = 2$ and $Y_\chi = 0$. The model files are implemented with these values.

The Yukawa interactions of the model and the vector-like mass term can be expressed as

$$-\mathcal{L}_Y = -\mathcal{L}_Y^{\mathrm{SM}} + \lambda^L\, \overline{Q^c}\, L\, \phi^* + \lambda^R\, \overline{u^c_R}\, e_R\, \phi^* + \frac{1}{2} M_\chi\, \overline{\chi}\, \chi + g_{S\chi} S\, \overline{\chi}\, \chi + \text{h.c.}\,, \tag{1.32}$$

whereas the scalar potential is given by

$$
\begin{aligned}
V = V_{\mathrm{SM}} &+ M_\phi^2|\phi|^2 + \frac{1}{2}M_S^2S^2 + \frac{1}{3}\lambda_{S_3}S^3 + \frac{1}{4}\lambda_{S_4}S^4 + \frac{1}{2}\lambda_\phi|\phi|^4 \\
&+ g_{SH}S^2|H|^2 + g_{H\phi}|\phi|^2|H|^2 + g_{S\phi}S^2|\phi|^2 + \kappa_{S\phi}S|\phi|^2 + \kappa_{SH}S|H|^2\,.
\end{aligned}
\tag{1.33}
$$

The Lagrangian in Eq. (1.32) is not generic. One could add more renormalizable operators involving the leptoquark, like $\overline{d^c}_R u_R \phi$ for instance. An operator of this kind, together with those in Eq. (1.32), would lead to rapid proton decay, thus it must be forbidden [47]. This can be achieved by imposing a discrete symmetry (for instance another $\mathbb{Z}_2$ under which the leptoquark and the SM quarks are odd and everything else is even). In the model file we simply omit these dangerous operators. The potential in Eq. (1.33), on the other hand, is generic. It could be simplified by introducing yet another $\mathbb{Z}_2$ to avoid terms with odd powers of S, and we would still have the necessary ingredients to fit the diphoton excess.

We assume the EWSB proceeds by the usual Higgs VEV in addition to the singlet VEV v_S, where the scalar singlet is expressed as

$$
S = v_S + \sigma_S\,. \tag{1.34}
$$

1.2.5 Two scalar leptoquark model

- **Reference:** [8]
- **Model name:** `Leptoquarks/TwoScalarLeptoquarks`

This model includes four additional scalars, two of which, Φ and Ω are leptoquarks. The other two are a singlet S, which is the 750 GeV resonance, and a scalar Θ which carries hypercharge. The additional quantum numbers for these new field are shown in Table 8.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$
S	1	$\mathbf{1}$	$\mathbf{1}$	0
Θ	1	$\mathbf{1}$	$\mathbf{1}$	1
Φ	1	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\frac{4}{3}$
Ω	1	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\frac{1}{3}$

Table 8. Additional scalar particles beyond the SM appearing in the two scalar leptoquark model.

The Yukawa interactions consistent with these new fields is given by

$$
-\mathcal{L}_Y = \mathcal{L}_Y^{\mathrm{SM}} + Y_\Theta L^T i\sigma_2 L\Theta + Y_\Omega Q_L^T i\sigma_2 L\Omega + Y_\Phi e_R^T d_R \Phi + \mathrm{h.c.}\,, \tag{1.35}
$$

where the generation indices have been suppressed. These indices are important so that one obtains non-zero contractions for the first term in the above Yukawa. The scalar potential of the model is given by

$$V = V_{\rm SM} + \sum_i \mu_\phi^2 |\phi_i|^2 + \frac{1}{2}\sum_{i,j} \lambda_{ij} |\phi_i|^2 |\phi_j|^2 + \left\{ \lambda S\,\Theta\,\Phi^\dagger \Omega + {\rm h.c.} \right\} , \tag{1.36}$$

where the index $i = \{S, \Theta, \Phi, \Omega\}$ runs over the introduced scalar fields. The potential has an additional $U(1)$ symmetry which can be gauged or global. In the implementation in `SARAH` we take it as global. This symmetry can be used to forbid terms such as $u_R d_R \Omega^\dagger$, which could potentially lead to proton decay.

One can generate neutrino masses at the two-loop level using a combination of the new Yukawa couplings and the VEV of the singlet. Beyond the neutral components of the Higgs doublet, the singlet also obtains a VEV of the form

$$S = v_S + \phi_S \,. \tag{1.37}$$

1.2.6 Georgi-Machacek model

- **Reference:** [9, 10]
- **Model name:** `Georgi-Machacek`

The model originally proposed in Ref. [48] extends the SM by one real scalar $SU(2)_L$-triplet η with $Y = 0$ and one complex scalar $SU(2)_L$-triplet with χ with $Y = 1$, which can be written as

$$\eta = \frac{1}{\sqrt{2}} \begin{pmatrix} \eta^0 & -\sqrt{2}\,(\eta^-)^* \\ -\sqrt{2}\eta^- & -\eta^0 \end{pmatrix} , \quad \chi = \frac{1}{\sqrt{2}} \begin{pmatrix} \chi^- & \sqrt{2}(\chi^0)^* \\ -\sqrt{2}\chi^{--} & -\chi^- \end{pmatrix} . \tag{1.38}$$

The CP-even components of the new scalars mix with the neutral SM Higgs. Usually, the lightest state is the 125 GeV Higgs boson, while the second mass eigenstate is identified with the 750 GeV resonance. The most compact form to write the Lagrangian in a $SU(2)_L \times SU(2)_R$ invariant form is to express the Higgs doublet and the two scalar triplets as

$$\Phi = \begin{pmatrix} \phi^{0*} & \phi^+ \\ \phi^- & \phi^0 \end{pmatrix} , \qquad \Delta = \begin{pmatrix} \chi^{0*} & \eta^+ & \chi^{++} \\ \chi^- & \eta^0 & \chi^+ \\ \chi^{--} & \eta^- & \chi^0 \end{pmatrix} . \tag{1.39}$$

In this form, the scalar potential reads

$$\begin{aligned} V(\Phi,\Delta) = {} & \mu_2^2 {\rm Tr}\Phi^\dagger\Phi + \frac{\mu_3^2}{2}{\rm Tr}\Delta^\dagger\Delta + \lambda_1 \left[{\rm Tr}\Phi^\dagger\Phi\right]^2 + \lambda_2 {\rm Tr}\Phi^\dagger\Phi\,{\rm Tr}\Delta^\dagger\Delta \\ & + \lambda_3 {\rm Tr}\Delta^\dagger\Delta\Delta^\dagger\Delta + \lambda_4 \left[{\rm Tr}\Delta^\dagger\Delta\right]^2 - \lambda_5 {\rm Tr}\left(\Phi^\dagger\sigma^a\Phi\sigma^b\right)\,{\rm Tr}\left(\Delta^\dagger t^a \Delta t^b\right) \\ & - M_1 {\rm Tr}\left(\Phi^\dagger\tau^a\Phi\tau^b\right)(U\Delta U^\dagger)_{ab} - M_2 {\rm Tr}\left(\Delta^\dagger t^a \Delta t^b\right)(U\Delta U^\dagger)_{ab} \,, \end{aligned}$$

τ^a and t^a are the $SU(2)$ generators for the doublet and triplet representations respectively, while U is given for instance in Ref. [49]. Because of the custodial symmetry, the VEVs for the triplets are identical, and there are no tree-level contributions to electroweak precision observables. The compact form of the potential can not be implemented in `SARAH`, so we have translated into the form of Eq. (1.38). For example:

$$\lambda_2 \mathrm{Tr}\Phi^\dagger\Phi\,\mathrm{Tr}\Delta^\dagger\Delta \quad \to \quad 4\lambda_2^a |H|^2 \mathrm{Tr}(\chi^\dagger\chi) + 2\lambda_2^b |H|^2 \mathrm{Tr}(\eta^\dagger\eta)\,. \tag{1.40}$$

This introduces more couplings, which have to be identical to preserve the custodial symmetry. A substitution rule to apply these relations is included in the model file.

The triplets receive VEV as

$$\langle\eta\rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v_\eta & 0 \\ 0 & -v_\eta \end{pmatrix}, \quad \langle\chi\rangle = \begin{pmatrix} 0 & v_\chi \\ 0 & 0 \end{pmatrix}, \tag{1.41}$$

which fulfil $v_\eta = v_\chi$ if the custodial symmetry is preserved.

1.3 Two-Higgs doublet models

Two-Higgs-doublet models (THDMs) are among the most common candidates to explain the diphoton excess. In this family of models, the 750 GeV resonance is typically a heavy Higgs, whose diphoton rate is enhanced with the addition of new states that contribute to its production cross-section and/or its decay width into pairs of photons. Several realisations of this idea can be found in the recent literature. Here we review a few representative examples.

THDM generalities

We first comment on some general features and conventions used in THDMs. For a complete review we refer to Ref. [50].

Couplings to fermions

We will consider three types of THDM, depending on the way the two Higgs doublets H_1 and H_2 couple to the SM fermions, as shown in this table:

	u_R	d_R	e_R
Type-I	H_2	H_2	H_2
Type-II	H_2	H_1	H_1
Type-III	both	both	both

We note that, by convention, the right-handed up-type quarks u_R always couple to H_2, and that the MSSM can be seen as a particular example of a type-II THDM. The type-I and type-II THDMs are usually said to be models with *natural flavour conservation*. This is because, in contrast to the general type-III THDM, flavour changing neutral currents are completely absent at tree-level, thus satisfying flavour constraints more easily. This is achieved by coupling each fermion species to only one Higgs doublet, which can be enforced using a discrete symmetry. For example, the type-I THDM couplings to fermions can be obtained by imposing that $H_1 \to -H_1$ is a symmetry of the Lagrangian. In our `SARAH` implementations we will not introduce these discrete symmetries explicitly, but simply allow only for those couplings that characterise the type of THDM under consideration. We point out that `SARAH` includes also other THDM versions, such as the lepton-specific or flipped ones. These versions can be easily combined with the extensions presented here for type I–III to include additional matter in order to explain the diphoton excess. Finally, we have also not implemented explicitly all representations for the vector-like states proposed in the literature so far. For instance, Ref. [51] introduces quarks with a very large hypercharge. However, it is easy to change the considered quantum numbers by changing the model files accordingly.

Scalar potential

We assume the following scalar potential for all THDMs considered below,

$$\begin{aligned} V = & \, m_{11}^2 |H_1|^2 + m_{22}^2 |H_2|^2 + \left[m_{12}^2 \, H_1^\dagger H_2 + \text{h.c.} \right] \\ & + \frac{\lambda_1}{2} |H_1|^4 + \frac{\lambda_2}{2} |H_2|^4 + \lambda_3 |H_1|^2 |H_2|^2 \\ & + \lambda_4 \left(H_1^\dagger H_2 \right) \left(H_2^\dagger H_1 \right) + \left[\frac{\lambda_5}{2} \left(H_1^\dagger H_2 \right)^2 + \text{h.c.} \right] . \end{aligned} \tag{1.42}$$

In principle, the quartic terms $\lambda_6 \left(H_1^\dagger H_1 \right) \left(H_1^\dagger H_2 \right)$ and $\lambda_7 \left(H_2^\dagger H_2 \right) \left(H_1^\dagger H_2 \right)$ are also allowed by the gauge symmetry. However, in our `SARAH` implementation we neglect these terms. This corresponds to assuming a global symmetry, which would be softly broken by the $m_{12}^2 \, H_1^\dagger H_2$ term.

Symmetry breaking pattern

The symmetry breaking pattern is the same in all THDM variations, namely

$$\langle H_i \rangle = \left\langle \begin{pmatrix} H_i^+ \\ H_i^0 \end{pmatrix} \right\rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_i \end{pmatrix} , \tag{1.43}$$

with $i = 1, 2$ running over the two Higgs-doublets.

In the following, we describe the THDM variations implemented in `SARAH` related to the diphoton excess. Broadly speaking, there are two main categories: THDMs with extra vector-like fermions and THDMs with extra scalar content. In the first case we have three examples, whereas for the second scenario we have only a single example.

1.3.1 Minimal vector-like THDM

- **Reference:** [11]
- **Model name:**
  `THDM+VL/min-3`
  `THDM+VL/min-8`

Ref. [11] considers a type-III THDM extended with two new vector-like colored fermions, S and Q^V, aiming at a simultaneous explanation of the diphoton excess and of the CMS hint for the Higgs lepton flavor violating (LFV) decay $h \to \tau\mu$ [52]. The type-III THDM has already been shown in several papers to be able to accommodate this LFV signal, see e.g. [53, 54], but [11] is the first work that addresses the diphoton excess at the same time. The representations of the new vector-like fermions under SM gauge group are shown in Table 9.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$
S	1	R_c	$\mathbf{1}$	Q
Q^V	1	R_c	$\mathbf{2}$	$Q+\frac{1}{2}$

Table 9. Additional fermion field content for the minimal vector-like THDM.

We choose the explicit realizations $R_c = 3$, $Q = 2$ (`THDM+VL/minTHDM-3`) and $R_c = 8$, $Q = 2$ (`THDM+VL/minTHDM-8`). In both cases, the additional interaction terms beyond those in the standard THDM are

$$-\mathcal{L}_{\text{new}} = M_Q^V \overline{Q_R^V} Q_L^V + M_S \overline{S_R} S_L + \lambda_i^D \overline{S_R} \tilde{H}_i Q_L^V + \lambda_i^S \overline{Q_R^V} H_i S_L + \text{h.c.}\,, \tag{1.44}$$

where M_Q^V and M_S are masses for the $SU(2)_L$-doublet and singlet vector-like fermions, respectively, and $\lambda_i^{D,S}$ (with $i = 1, 2$) are new Yukawa interactions.

1.3.2 THDM with exotic vector-like fermions

- **References:** [12, 13]
- **Model name:**
  `THDM+VL/Type-I-VL`
  `THDM+VL/Type-II-VL`

The THDMs in Refs. [12, 13] consider a less minimal framework where three generations of new vector-like fermions are added to the standard THDM scenario. Both type-I and type-II THDMs are studied. Furthermore, the vector-like leptons have exotic hypercharge values shown in Table 10.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$\mathbb{Z}_2$
Q^V	3	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$	$-$
d^V	3	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	$-$
u^V	3	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	$-$
L^V	3	$\mathbf{1}$	$\mathbf{2}$	$-\frac{3}{2}$	$-$
e^V	3	$\mathbf{1}$	$\mathbf{1}$	-2	$-$
ν^V	3	$\mathbf{1}$	$\mathbf{1}$	-1	$-$

Table 10. Additional fermion field content for the THDM with exotic vector-like fermions. All SM fields are even under the imposed $\mathbb{Z}_2$ discrete symmetry.

The vector-like lepton states can be decomposed (or denoted) as

$$L^V = \begin{pmatrix} \ell'^- \\ \ell^{--} \end{pmatrix}, \qquad e^V \equiv \left(e^V\right)^{--}, \qquad \nu^V \equiv \left(\nu^V\right)^{-}, \tag{1.45}$$

where we explicitly highlight the presence of doubly-charged leptons. The additional interaction terms beyond those in the standard THDMs are

$$\begin{aligned} -\mathcal{L}_{\text{new}} = {} & M_Q^V \overline{Q_R^V} Q_L^V + M_L^V \overline{L_R} L_L + M_d^V \overline{d_R^V} d_L^V + M_u^V \overline{u_R^V} u_L^V + M_\nu^V \overline{\nu_R^V} \nu_L^V + M_e^V \overline{e_R^V} e_L^V \\ & + Y_{dL}^V \overline{d_R^V} \tilde{H}_1 Q_L^V + Y_{dR}^V \overline{Q_R^V} H_1 d_L^V + Y_{uL}^V \overline{u_R^V} H_2 Q_L^V + Y_{uR}^V \overline{Q_R^V} \tilde{H}_2 u_L^V \\ & + Y_{eL}^V \overline{e_R^V} \tilde{H}_1 L_L^V + Y_{eR}^V \overline{L_R^V} H_1 e_L^V + Y_{\nu L}^V \overline{\nu_R^V} H_2 L_L^V + Y_{\nu R}^V \overline{L_R^V} \tilde{H}_2 \nu_L^V + \text{h.c.}\,. \end{aligned} \tag{1.46}$$

Here M_i^V are the bare vector-like masses and Y_i^V are Yukawa couplings. For simplicity, we assume that the vector-like states do not mix with the SM fermions. This can be enforced with a discrete $\mathbb{Z}_2$ symmetry under which the vector-like states are odd and the rest of the states even.

1.3.3 THDM with SM-like vector-like fermions

- **Reference:** [14]
- **Model name:**
  `THDM+VL/Type-I-SM-like-VL`
  `THDM+VL/Type-II-SM-like-VL`

The THDM proposed in [14] introduces vector-like copies of the SM quarks and leptons, the $SU(2)_L$ doublets Q^V and L^V, as well as the $SU(2)_L$ singlets d^V, u^V, ν^V and e^V with their respective quantum numbers shown in Table 11.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$\mathbb{Z}_2$
Q^V	1	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$	$-$
d^V	1	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	$-$
u^V	1	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	$-$
L^V	1	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	$-$
e^V	1	$\mathbf{1}$	$\mathbf{1}$	-1	$-$
ν^V	1	$\mathbf{1}$	$\mathbf{1}$	0	$-$

Table 11. Additional fermion field content for the THDM with SM-like vector-like fermions. All SM fields are even under the imposed $\mathbb{Z}_2$ discrete symmetry.

In contrast to other THDMs proposed to explain the diphoton excess, where low values of $\tan\beta$ are taken in order to increase the heavy Higgs coupling to the top quark and induce a large gluon fusion cross-section, this paper considers moderate and large values of $\tan\beta$. In this case, the heavy Higgs production is mainly induced by the new vector-like colored states. Moreover, the advantage of using largish $\tan\beta$ values is that the total decay width of the heavy Higgs is suppressed, thus allowing for a smaller $\Gamma(H \to \gamma\gamma)$ to explain the excess.

The new interaction terms take the same form as those in Eq. (1.46). Furthermore, as in the previous models, we assume that the new vector-like states do not mix with the SM fermions. For this reason, we introduce a $\mathbb{Z}_2$ parity under which the vector-like states are odd and the rest of the states even. Finally, Ref. [14] considers two variants of this scenario in what concerns the couplings to the SM fermions: a type-I THDM and a type-II THDM.

1.3.4 THDM with a real scalar septuplet

- **Reference:** [15, 16]
- **Model name:** `THDM/ScalarSeptuplet`

Ref. [15] explores an extension of the type-III THDM with new scalar $SU(2)_L$ representations, namely an inert complex Higgs triplet and a real scalar septuplet. A more general analysis can be found in Ref. [16], where a generic real scalar $SU(2)_L$ multiplet is considered. In both papers, the 750 GeV resonance is identified with the usual heavy Higgs of the THDM, with the additional scalars being introduced to push the diphoton rate to higher values.

For the realization of this idea here we focus on the septuplet case, closely following Ref. [15]. The introduction of such large scalar multiplets has become fairly popular recently in the context of minimal DM scenarios [55], although [15] does not explore any dark matter implications. It only takes advantage of the multicharged states contained in the septuplet which, due to their couplings to the THDM doublets, lead to a large diphoton rate for the heavy Higgs H.

The type-III THDM is extended with the addition of a real scalar which transforms as a **7** under $SU(2)_L$, see Table 12.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$
T	1	**1**	**7**	0

Table 12. Additional fermion field content for the scalar septuplet extended THDM.

It proves convenient to use tensor notation for the septuplet. This is the usual choice, see Refs. [56, 57], and the one implemented in SARAH. The septuplet T is represented by a symmetric tensor with six indices, T^{ijklmn}, all of which can take values of either 1 or 2. The relation with the vector notation, employed for example in [15], is given by

$$T \equiv \begin{pmatrix} T^{+++} \\ T^{++} \\ T^{+} \\ T^{0} \\ T^{-} \\ T^{--} \\ T^{---} \end{pmatrix} = i \begin{pmatrix} +T^{111111} \\ +\sqrt{6}\,T^{111112} \\ +\sqrt{15}\,T^{111122} \\ -\sqrt{20}\,T^{111222} \\ -\sqrt{15}\,T^{112222} \\ +\sqrt{6}\,T^{122222} \\ -T^{222222} \end{pmatrix} . \tag{1.47}$$

The prefactor i and the sign for each component are introduced in order to satisfy $T^c = T$ where T^c denotes charge conjugation of the field T, and T^0 is a real scalar. The new potential terms involving the septuplet T are

$$V_T = M_T^2\, T^2 + \sum_{i=1}^{2} \lambda_i^T \, \left[T^4\right]_i + \lambda_3^T \, |H_1|^2 \, T^2 + \lambda_4^T \, |H_2|^2 \, T^2 \, . \tag{1.48}$$

We note that two independent gauge invariant contractions are possible in case of T^4, which is reflected by the second term of Eq. (1.48). Finally, the authors of [15] assume a minimum of the scalar potential with $\langle T \rangle = 0$. In this case, the components of the septuplet do not mix with the scalar doublets, but only participate in the $H \to \gamma\gamma$ rate due to the interaction terms $\lambda^T_{3,4}$.

2 $U(1)$ extensions of the SM

In this section we consider a class of models which extend the SM by a new $U(1)_X$ gauge group. One typically introduces, beyond the SM Higgs doublet, new scalars charged under $U(1)_X$, which serve two purposes: (i) a linear combination of them results in the 750 GeV particle, and (ii) via spontaneous symmetry breaking they give mass to the new gauge boson, the Z' boson. Typically, one also introduces new fermions charged under the $U(1)_X$ symmetry that can either be singlets under the SM gauge group, hence forming a dark or hidden sector, or vector-like under the SM. An advantage of these models is that, through a suitable choice of charge assignments under $U(1)_X$, one can avoid flavour constraints present when allowing the additional quarks to decay. Finally, the presence of a massive Z' boson can lead to new collider signals, well studied in the literature, which can serve as a smoking gun for these types of models. We note that the mixing matrix for the neutral gauge bosons can be parametrised by two angles, Θ and Θ', with Θ' highly constrained by LEP data.

We distinguish two cases in the following: models in which the SM fermions are charged under the new Abelian group, and models in which they are singlet.

2.1 Models with SM states uncharged under the new $U(1)$

2.1.1 Dark $U(1)'$ extension

- **Reference:** [17]
- **Model name:** `U1Extensions/darkU1`

This model is based on a gauged $U(1)_X$ extension of the SM with a dark sector that includes three generations of dark $SU(2)_L$-singlet fermions and a DM scalar candidate.

The properties of the new particles introduced in this model are described in Table 13. We have added three right-handed neutrinos ν^i_R, neutral under $U(1)_X$, to the original model, since N_R (their dark partners) were considered here. Φ is the scalar field responsible for the spontaneous symmetry breaking of $U(1)_X$, while X is the DM candidate. The $U(1)_X$ charges a, b are left arbitrary, with their assignment chosen such that they fulfil the anomaly cancellation conditions.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_X$
Φ	1	$\mathbf{1}$	$\mathbf{1}$	0	$a+b$
X	1	$\mathbf{1}$	$\mathbf{1}$	0	a
ν_R	3	$\mathbf{1}$	$\mathbf{1}$	0	0
E_L	3	$\mathbf{1}$	$\mathbf{1}$	-1	a
E_R	3	$\mathbf{1}$	$\mathbf{1}$	-1	$-b$
N_L	3	$\mathbf{1}$	$\mathbf{1}$	0	$-a$
N_R	3	$\mathbf{1}$	$\mathbf{1}$	0	b
U_L	3	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	$-a$
U_R	3	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	b
D_L	3	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	a
D_R	3	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	$-b$

Table 13. New fermions and scalar fields of the `darkU1` and their charge assignments under the gauge symmetry $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_X$. The scalar and fermion fields are shown in the top and bottom of the table respectively.

The Yukawa interactions and the scalar potential including new fields in the dark sector are described by

$$\begin{aligned} -\mathcal{L}_Y^{\text{new}} &= Y'_E\, \overline{E_R}\, \Phi^*\, E_L + Y'_N\, \overline{N_R}\, \Phi\, N_L + Y'_U\, \overline{U_R}\, \Phi\, U_L + Y'_D\, \overline{D_R}\, \Phi^*\, D_L + Y_{XE}\, \overline{e_R} X^*\, E_L \\ &\quad + Y_{XU}\, \overline{u_R}\, X\, U_L + Y_{XD}\, \overline{d_R}\, X^*\, D_L + Y_{XN}\, \overline{\nu_R}\, X\, N_L + \text{h.c.}\,, \qquad (2.1) \\ V &= \mu^2 |H|^2 + \lambda |H|^4 + \mu_\Phi^2 |\Phi|^2 + \mu_X^2 |X|^2 + \lambda_\Phi |\Phi|^4 + \lambda_X |X|^4 + \lambda_{H\Phi} |H|^2 |\Phi|^2 \\ &\quad + \lambda_{HX} |H|^2 |X|^2 + \lambda_{X\Phi} |X|^2 |\Phi|^2\,, \qquad (2.2) \end{aligned}$$

where H denotes the SM Higgs doublet. For $a = b = 1$, an extra term $\Phi^\dagger X^2$ would be allowed in the potential, which breaks $U(1)_X$ down to a Z_2 subgroup after Φ develops a non-zero VEV. Likewise, for $3a = (a+b)$, there appears an extra term $\Phi^\dagger X^3$, which breaks $U(1)_X$ down to a Z_3 subgroup after Φ develops a nonzero VEV. These possibilities are not considered here.

The gauge symmetry is broken after H and Φ get non-zero VEVs, v and v_Φ respectively, while X does not receive a VEV[2]. The scalar fields after EWSB can be expressed as

$$H = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v + \phi_H + i\sigma_H \end{pmatrix}\,, \qquad \Phi = \frac{1}{\sqrt{2}} \left(v_\Phi + \phi_\Phi + i\sigma_\Phi \right)\,, \qquad (2.3)$$

[2] We also checked that this condition remains valid at NLO.

where ϕ_i and σ_i are the CP-even and odd components respectively. The 750 GeV candidate is the mixture of the SM Higgs ϕ_H and ϕ_Φ, which leads to constraints on $\lambda_{H\Phi}$. It is produced in gluon-gluon fusion via loops of the dark quarks, and it decays into diphotons via loops comprised only of charged dark fermions, assuming that the mixing with the SM-like Higgs is negligibly small.

2.1.2 Hidden $U(1)$

- **Reference:** [18]
- **Model name:** `U1Extensions/hiddenU1`

This is a particularly simple realisation of a gauged $U(1)_X$ extension of the SM. As previously, a hidden $U(1)_X$ is added to the SM gauge group, under which all SM particles are singlets. A scalar S_1 is added for its spontaneous symmetry breaking, and a further scalar S_2 is added which is a singlet under the entire gauge symmetry of the model. Here we assume that both S_1 and S_2 can develop a VEV, in principle.

The 750 GeV candidate is considered to be predominantly composed by S_2. To explain the diphoton signal, a vector-like quark is also included, carrying the same charge under $U(1)_X$ as the S_1 field. Hence, it couples only to S_2 due to the choice of charge assignments.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_X$
S_1	1	$\mathbf{1}$	$\mathbf{1}$	0	a
S_2	1	$\mathbf{1}$	$\mathbf{1}$	0	0
X_L	1	$\mathbf{3}$	$\mathbf{1}$	Y_X	a
X_R	1	$\mathbf{3}$	$\mathbf{1}$	Y_X	a

Table 14. New fermions and scalar fields of the `hiddenU1` and their charge assignments under the gauge symmetry $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_X$. The scalar and fermion fields are shown in the top and bottom of the table respectively.

In Table 14 the hypercharge of the new vector-like quark Y_X is left arbitrary. The implemented case is the most favourable one in terms of the diphoton excess, $Y_X = 2/3$, which allows a mixing with the up-type quarks. We did not considered the case of adding also a vector-like lepton to the spectrum, which would lead to even larger rates. The $U(1)_X$ charge of the vector-like quark fields a does not affect the diphoton rate, but it has to be the same as for S_1, which is relevant for the Z' boson mass.

The scalar potential, with the usual doublet Higgs field H is given by

$$\begin{aligned} V = &\mu_H^2|H|^2 + \mu_{S_1}^2|S_1|^2 + \mu_{S_2}^2 S_2^2 - \lambda_H|H|^4 - \lambda_{HS_1}|H|^2|S_1|^2 - \lambda_{S_1}|S_1|^4 \\ &- \lambda_{S_2}S_2^4 - \lambda_{HS_2}|H|^2S_2^2 - \lambda_{S_1S_2}|S_1|^2S_2^2 - \sigma_1 S_2^3 - \sigma_2|H|^2S_2 - \sigma_3|S_1|^2S_2 \,. \end{aligned} \tag{2.4}$$

This potential leads to a mixing between the three physical neutral scalars. The structure of the new scalar fields is given by

$$S_1 = \frac{1}{\sqrt{2}}\left(v_{S_1} + \phi_{S_1} + i\sigma_{S_1}\right)\,, \qquad S_2 = v_{S_2} + \phi_{S_2}\,, \tag{2.5}$$

where once again ϕ_i and σ_i are the CP-even and odd components respectively. As discussed in [1] we have allowed all scalar fields to obtain VEVs.

The Yukawa interactions and fermionic mass terms of the hidden sector read

$$-\mathcal{L}^{\text{new}} = M_X\,\overline{X_R}\,X_L + Y_{X_L}\,\overline{u_R}\,S_1^*\,X_L + f_X\,\overline{X_R}\,S_2X_L \;+\text{h.c.}\,. \tag{2.6}$$

The mixing of the vector-like quark X with SM quarks via the interaction with S_1 is kept small, and purely serves the purpose of letting the new quark decay.

2.1.3 Simple $U(1)$

- **Reference:** [19]
- **Model name:** `U1Extensions/simpleU1`

This $U(1)_X$ extension of the SM augments its particle content by a pair of exotic vector-like quarks, χ_1 and χ_2, doublets under $SU(2)_L$ and with hypercharge 7/6, and by one complex scalar Σ responsible for the spontaneous breaking of $U(1)_X$. The $U(1)_X$-breaking Higgs boson is the 750 GeV candidate. The particle content beyond the SM is summarized in Table 15.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_X$
Σ	1	**1**	**1**	0	1
χ_1	1	**3**	**2**	7/6	1
χ_2	1	**3**	**2**	7/6	2

Table 15. Particle content of the `simpleU1` beyond the SM fields. The scalar and fermion fields are shown in the top and bottom of the table respectively. The exotic χ_1 and χ_2 quarks are vector-like.

The scalar potential of the model is given by

$$V = -\mu_H^2 |H|^2 - \mu_\Sigma^2 |\Sigma|^2 + \lambda_H |H|^2 + \lambda_{H\Sigma} |H|^2 |\Sigma|^2 + \lambda_\Sigma |\Sigma|^4 \,, \tag{2.7}$$

while the Yukawa and fermionic mass terms read

$$\mathcal{L} = \mathcal{L}_Y^{\rm SM} - M_1 \, \overline{\chi_{1R}} \, \chi_{1L} - M_2 \, \overline{\chi_{2R}} \, \chi_{2L} - \lambda_1 \, \Sigma \, \overline{\chi_{2R}} \, \chi_{1L} - \lambda_2 \, \Sigma^* \overline{\chi_{1R}} \chi_{2L} \ . \tag{2.8}$$

Note that the original model proposed in [19] contains an effective operator

$$\mathcal{L} \supset -\frac{1}{\Lambda} \Sigma^* H \chi_1 u^c \,.$$

As `SARAH` cannot handle effective operators this term is dropped from the model and subsequent constraints pertaining to stable charged particles are ignored. Expanding the new scalar field yields

$$\Sigma = \frac{1}{\sqrt{2}} \left(v_\Sigma + \phi_\Sigma + i \sigma_\Sigma \right) \,, \tag{2.9}$$

where once again ϕ_i and σ_i are the CP-even and odd components respectively.

2.1.4 Scotogenic $U(1)$ Model

- **Reference:** [20]
- **Model name:** `U1Extensions/scotoU1`

The matter particle content of the Scotogenic $U(1)$ Model is summarized in Table 16, where, in addition to the fields charged under $U(1)_D$, we introduce three copies of right-handed neutrinos ν_R which are singlets under the full gauge group. Note, that the $U(1)_D$ charge of the H' field has been changed to -1 compared to Ref. [20] in order to make the Yukawa interaction terms gauge invariant. The SM fields are not charged under the new $U(1)$ gauge group. The discrete $\mathbb{Z}_2$ symmetry is introduced to stabilize the dark matter candidates N.

The scalar potential reads

$$\begin{aligned} V = & -\mu_1^2 |H|^2 + \mu_2^2 |H'|^2 - \mu_s^2 |\Phi|^2 + \lambda_1 |H|^4 + \lambda_2 |H'|^4 + \lambda_s |\Phi|^4 \\ & + \lambda_h |H|^2 |\Phi|^2 + \lambda_{h'} |H'|^2 |\Phi|^2 + \lambda_3 |H|^2 |H'|^2 + \lambda_4 |H^\dagger H'|^2 \,, \end{aligned} \tag{2.10}$$

The term $\frac{\lambda_5}{2}[(H^\dagger H')^2 + \text{h.c.}]$ proposed in Ref. [20] has been omitted here, because it is not invariant under $U(1)_D$ gauge transformations.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_D$	$\mathbb{Z}_2$
Φ	1	**1**	**1**	0	2	+
H'	1	**1**	**2**	$\frac{1}{2}$	-1	$-$
ν_R	3	**1**	**1**	0	0	+
T	1	**3**	**1**	$\frac{2}{3}$	-1	$-$
T'	1	**3**	**1**	$\frac{2}{3}$	1	$-$
N	1	**1**	**1**	0	-1	$-$

Table 16. Matter particle content of the `scotoU1` beyond the SM fields. The scalar and fermion fields are shown in the top and bottom of the table, respectively. The fermions T, T', N are vector-like degrees of freedom (4-component spinors).

The Yukawa interactions beyond the SM read

$$\begin{aligned}\mathcal{L} \supset &+Y_\nu \nu_R H \ell + Y_T \overline{T_R} H' q + y_N \overline{N_R} H' L + m_T(\overline{T}T + \overline{T'}T') \\ &+ M_D \overline{N} N + \eta_1 \overline{T'_R} \Phi T_L + \eta_2 \overline{T_R} \Phi^* T'_L + \eta_3 \overline{N^c_R} \Phi N_L + \text{h.c.}\,,\end{aligned} \tag{2.11}$$

where $H = (H^+, H^0)$ is the SM Higgs field and q and ℓ the SM left-handed quark and lepton doublets.

The $U(1)_D$ symmetry is eventually broken by the VEV of the scalar field Φ which one can decompose as

$$\langle \Phi \rangle = \frac{1}{\sqrt{2}}(v_S + S + iA_S)\,, \tag{2.12}$$

whereas the CP-even component S is considered by the authors as the candidate for the 750 GeV resonance. The spontaneous symmetry breaking leaves the $\mathbb{Z}_2$ parity intact, and H' will therefore not develop any VEV.

2.2 Models with SM states charged under the new $U(1)$

2.2.1 $U(1)_{B-L}$ model with unconventional $B-L$ charges

- **Reference:** [21]
- **Model name:** `U1Extensions/BL-VL`

This model is based on Refs. [58, 59]. It considers a gauged $U(1)_{B-L}$ extension of the SM with an unconventional $B-L$ charge assignment for the right-handed neutrinos ν_R, and further requires 3 extra Dirac neutrinos N. It was originally proposed to explain the smallness of neutrino masses if neutrinos were Dirac particles. Finally, it

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_{B-L}$
Φ	1	**1**	**2**	$\frac{1}{2}$	0
χ_2	1	**1**	**1**	0	2
χ_3	1	**1**	**1**	0	3
χ_6	1	**1**	**1**	0	-6
Q_L	3	**3**	**2**	$\frac{1}{6}$	$\frac{1}{3}$
u_R	3	**3**	**1**	$\frac{2}{3}$	$\frac{1}{3}$
d_R	3	**3**	**1**	$-\frac{1}{3}$	$\frac{1}{3}$
L_L	3	**1**	**2**	$-\frac{1}{2}$	-1
l_R	3	**1**	**1**	-1	-1
ν_R^1	1	**1**	**1**	0	5
ν_R^2	1	**1**	**1**	0	-4
ν_R^3	1	**1**	**1**	0	-4
N_L	3	**1**	**1**	0	-1
N_R	3	**1**	**1**	0	-1
X_L	1	**3**	**1**	$\frac{2}{3}$	3
X_R	1	**3**	**1**	$\frac{2}{3}$	0
Y_L	1	**3**	**1**	$-\frac{2}{3}$	-3
Y_R	1	**3**	**1**	$-\frac{2}{3}$	0

Table 17. Particle content and charge assignments of the `BL-VL` model. The scalar and fermion fields are shown in the top and bottom of the table respectively.

also contains a scalar DM candidate, stabilised by the residual $\mathbb{Z}_3$ symmetry from the breaking of the $B-L$ symmetry. To fit the diphoton excess, new vector-like quarks are added so that the 750 GeV scalar, a linear combination of the SM Higgs and the χ fields, can be produced in gluon-gluon fusion. The particle content and the quantum numbers for this model are summarised in Table 17. In addition to the SM particle content, the model features three right-handed neutrinos ν_R^i, three pairs of $SU(2)_L$ singlet heavy fermions $N^i_{L,R}$, and two pairs of exotic quarks $X_{L,R}$, $Y_{L,R}$ which carry color and electromagnetic charges but are singlets under $SU(2)_L$.

The scalar potential of the model reads

$$\begin{aligned}V = &-\mu_0^2|\Phi|^2 + m_2^2|\chi_2|^2 - \mu_3^2|\chi_3|^2 - \mu_6^2|\chi_6|^2 + \frac{1}{2}\lambda_0|\Phi|^4 + \frac{1}{2}\lambda_2|\chi_2|^4 + \frac{1}{2}\lambda_3|\chi_3|^4 + \frac{1}{2}\lambda_6|\chi_6|^4 \\ &+ \lambda_{02}|\chi_2|^2|\Phi|^2 + \lambda_{03}|\chi_3|^2|\Phi|^2 + \lambda_{06}|\chi_6|^2|\Phi|^2 + \lambda_{23}|\chi_2|^2|\chi_3|^2 + \lambda_{26}|\chi_2|^2|\chi_6|^2 \\ &+ \lambda_{36}|\chi_3|^2|\chi_6|^2 + \left[\frac{1}{2}f_{36}(\chi_3^2\chi_6) + \frac{1}{6}\lambda'_{26}(\chi_2^3\chi_6) + \text{h.c.}\right] . \end{aligned} \tag{2.13}$$

The residual global $\mathbb{Z}_3$ symmetry protects the singlet scalar χ_2 from acquiring a VEV, i.e. $\langle\chi_2\rangle = 0$. All leptons carry a charge $\omega = e^{2i\pi/3}$ under $\mathbb{Z}_3$. The CP-even degree of freedom of χ_2 is the DM candidate. The Yukawa interactions of the new sector read

$$\begin{aligned}
-\mathcal{L}_Y^{\text{new}} &= Y_{NR}\,\overline{N_R}\,\widetilde{H}\,L_L + f_X\,\overline{X_R}\,\chi_3^*\,X_L + f_Y\,\overline{Y_R}\,\chi_3\,Y_L \\
&+ f_N\,\overline{\nu_{R_{2,3}}}\,\chi_3^*\,N_L + f_{N6}\,\overline{\nu_{R_1}}\,\chi_6^*\,N_L \\
&+ f_{NL}\,\chi_2\,\overline{N_L^c}\,N_L + f_{NR}\,\chi_2\,\overline{N_R^c}\,N_R + m_N\,\overline{N_R}\,N_L + \text{h.c.}\,.
\end{aligned} \tag{2.14}$$

In the set up considered in Ref [21], the 750 GeV scalar is given by the combination $(\chi_6-\chi_3)$, that couples to gluons and photons via loops of X and Y fermions proportional to the Yukawa couplings f_X and f_Y. The fit to the diphoton signal was studied in a simplified scenario, where special relations are imposed to the scalar parameters. As already described in [1], these relations are not protected by symmetries, and hence lead to a large amounts of fine tuning.

2.2.2 Sample of $U(1)'$ models based on different charge assignments

- **Reference:** [22]
- **Model name:** `U1Extensions/VLsample`

The particle content of this model is shown in Table 18. Similarly to all the previous models, a complex scalar S is added to break the $U(1)_X$ symmetry. In this paper, the mixing between the scalars is kept small, hence the 750 GeV candidate is predominantly the CP-even part of the $U(1)_X$-breaking field. New doublet and singlet vector-like quarks, charged under $U(1)_X$, are added to fit the diphoton signal strength. In this model, the SM fermions also carry X-charges, which are fixed according to the anomaly cancellation conditions. Only the Higgs doublet is not charged, while the charge of the new scalar field S is double (in absolute magnitude) that of the vector-like quarks.

There are several possibilities to assign the $U(1)_X$ charges in an anomaly-free way, with different physical interpretations [22]:

- $X \equiv B-L$: $b=0,\ k=-1,\ m=1/3$
- $X \equiv B+L$: $b=-1,\ k=1,\ m=1/3$
- $X \equiv B$: $b=-1/2,\ k=0,\ m=1/3$
- $X \equiv L$: $b=-1/2,\ k=1,\ m=0$

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_X$
H	1	**1**	**2**	$-\frac{1}{2}$	0
S	1	**1**	**1**	0	$-2b$
Q_L	3	**3**	**2**	$\frac{1}{6}$	m
u_R	3	**3**	**1**	$\frac{2}{3}$	m
d_R	3	**3**	**1**	$-\frac{1}{3}$	m
L_L	3	**1**	**2**	$-\frac{1}{2}$	k
e_R	3	**1**	**1**	-1	k
ν_R	3	**1**	**1**	0	k
X_L	1	**3**	**2**	a	b
X_R	1	**3**	**2**	a	$-b$
y_{1L}	1	**3**	**1**	$a+\frac{1}{2}$	$-b$
y_{2L}	1	**3**	**1**	$a-\frac{1}{2}$	$-b$
y_{1R}	1	**3**	**1**	$a+\frac{1}{2}$	b
y_{2R}	1	**3**	**1**	$a-\frac{1}{2}$	b

Table 18. Particle content of the `VLsample`. The scalar and fermion fields are shown in the top and bottom of the table respectively.

Note that in all interpretations a remains a free parameter.

The Yukawa interactions of the extra particle content are given by

$$\begin{aligned}\mathcal{L}_Y^{\text{new}} &= Y_V^1\overline{y_{1R}}S^* y_{1L} + Y_V^2\overline{y_{2R}}S^* y_{2L} + Y_V^3\overline{X_R}SX_L \\ &+ \eta_1\overline{X_R}Hy_{1L} + \eta_2\overline{X_R}\tilde{H}y_{2L} + \eta_3\overline{y_{2R}}HX_L + \eta_4\overline{y_{1R}}\tilde{H}X_L \ + \text{h.c.}\end{aligned} \tag{2.15}$$

Finally, the scalar potential is given by

$$-V = \mu_H^2|H|^2 + \mu_S^2|S|^2 - \lambda_H|H|^4 - \lambda_S|S|^4 - \lambda_{HS}|S|^2|H|^2\,, \tag{2.16}$$

the symmetry breaking pattern being

$$\langle H\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} v \\ 0\end{pmatrix}\,, \quad \langle S\rangle = \frac{1}{\sqrt{2}}\,v_S\ , \tag{2.17}$$

while the expansion of the scalar fields in to CP-even and odd components is analogous to Eq. (2.3).

2.2.3 Model with flavour-nonuniversal quark $U(1)'$ charges

- **Reference:** [23]

- **Model name:** `U1Extensions/nonUniversalU1`

In this model the first generation of left-handed SM quarks carries a $U(1)_X$ charge while the second and third generations do not. In this way it is possible to add exotic quarks which are vector-like under the SM gauge group and achieve anomaly cancellation with less than three generations. The scalar sector then needs to be extended with a second Higgs doublet, with a different $U(1)_X$ charge compared to the first ones, in order to have Yukawa interactions for all quark families. A further complex scalar field S, which is a singlet under the SM gauge group but charged under $U(1)_X$, is also added for the spontaneous symmetry breaking of the $U(1)_X$ symmetry.

The charge assignments that cancel all anomalies are given in [60] and are summarised in Table 19.[3]

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_X$
H_1	1	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	$-\frac{2}{3}$
H_2	1	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	$-\frac{1}{3}$
S	1	$\mathbf{1}$	$\mathbf{1}$	0	$-\frac{1}{3}$
Q_L^1	1	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$	$\frac{1}{3}$
Q_L^i	2	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$	0
u_R	3	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	$\frac{2}{3}$
d_R	3	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	$-\frac{1}{3}$
L_L	3	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	$-\frac{1}{3}$
e_R	3	$\mathbf{1}$	$\mathbf{1}$	-1	-1
ν_R	3	$\mathbf{1}$	$\mathbf{1}$	0	$\frac{1}{3}$
T_L	1	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	$\frac{1}{3}$
T_R	1	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	$\frac{2}{3}$
J_L	2	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	0
J_R	2	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	$-\frac{1}{3}$

Table 19. Particle content of the `nonUniversalU1`. The scalar and fermion fields are shown in the top and bottom of the table respectively.

[3] In an updated version of their paper, the authors of Ref. [23] have added a further scalar σ with the same quantum numbers as S which shall be the dark matter candidate as well as extra fermionic singlets in order to allow for a seesaw mechanism in the neutrino sector.

The Yukawa Lagrangian of the model (ignoring here flavour indices) is given by

$$\begin{aligned}\mathcal{L}_Y = & \; h_1^D \overline{d_R} H_1 Q_L^1 + h_1^U \overline{u_R} \tilde{H}_1 Q_L^i + h_2^D \overline{d_R} H_2 Q_L^i + h_2^U \overline{u_R} \tilde{H}_2 Q_L^1 + h_1^J \overline{J_R} H_1 Q_L^1 \\ & + h_2^J \overline{J_R} H_2 Q_L^i + h_2^T \overline{T_R} \tilde{H}_2 Q_L^1 + h_1^T \overline{T_R} \tilde{H}_1 Q_L^i + h_X^U \overline{u_R} S^* T_L + h_X^T \overline{T_R} S^* T_L \\ & + h_X^D \overline{d_R} S J_L + h_X^J \overline{J_R} S J_L + Y_e \overline{e_R} H_1 L_L + Y_v \overline{\nu_R} \tilde{H}_1 L_L + \text{h.c.}\,, \end{aligned} \tag{2.18}$$

and the scalar potential is

$$\begin{aligned} V = & \; \mu_{11}^2 |H_1|^2 + \mu_{22}^2 |H_2|^2 + \mu_S^2 |S|^2 + \lambda_1 |H_1|^4 + \lambda_2 |H_2|^4 + \lambda_3 |H_1|^2 |H_2|^2 - \lambda_{H_1 S} |S|^2 |H_1|^2 \\ & - \lambda_S |S|^4 + \lambda_4 (H_2^\dagger H_1)(H_1^\dagger H_2) - \lambda_{H_2 S} |S|^2 |H_2|^2 + \left\{ \kappa_{HS} H_1^\dagger H_2 S + \text{h.c.} \right\} . \end{aligned} \tag{2.19}$$

The pattern of EWSB follows

$$H_i = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_i + \phi_i + i\sigma_i \end{pmatrix} , \qquad S = \frac{1}{\sqrt{2}} \left(v_S + \phi_S + i\sigma_S \right) , \tag{2.20}$$

where once again ϕ_i and σ_i are the CP-even and CP-odd components respectively. In order to obtain a massive pseudo-scalar state, κ_{HS} needs to be non-zero. Hence, for keeping the $S - H_i$ mixing small while making the pseudo-scalar massive, either the condition $v_S \gg v$ must hold, or κ_{HS} must be small, in conjunction with $\frac{v_1}{v_2} \to 0$ or ∞. ϕ_S, the CP-even component of S, is then identified with the 750 GeV candidate.

2.2.4 Leptophobic $U(1)$ model

- **Reference:** [24]
- **Model name:** `U1Extensions/U1Leptophobic`

This model is inspired by an E_6 Grand Unified Theory (GUT), but the authors only consider the low energy version where the SM gauge group is augmented by an extra gauged $U(1)_X$ symmetry. This extra $U(1)$ symmetry has zero charges for both left- and right-handed leptons making it entirely lepotophobic. However it is impossible to arrange for this to happen by taking linear combinations of the $U(1)_\chi$ and $U(1)_\psi$ groups that appear from the breakdown of E_6. Instead, the charges from this extra $U(1)$ can only be obtained by including gauge kinetic mixing, so that the introduced mixture of $U(1)_Y$ charges exactly cancel the non-zero leptonic charges. This can be done with the $U(1)_\eta$ gauge symmetry and the charges used in this model correspond exactly to those given in Table I of Ref. [61]. It is these charges rather than the charges of the $U(1)_\eta$ which are set in this model. Of course one may discard the E_6 motivation and treat it as an ad-hoc choice of $U(1)$ charges.

The model contains two Higgs doublets, a complex scalar SM singlet (Φ), charged under $U(1)_X$, plus right handed neutrinos and other new fermions, charged under both the SM and the $U(1)_X$ symmetries. The latter are odd under an imposed $\mathbb{Z}_2$symmetry, so that the lightest neutral fermion is a DM candidate.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_X$	$\mathbb{Z}_2$
H_1	1	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	0	$+$
H_2	1	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	-1	$+$
Φ	1	$\mathbf{1}$	$\mathbf{1}$	0	-1	$+$
Q_L	3	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$	$-\frac{1}{3}$	$+$
u_R	3	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	$\frac{2}{3}$	$+$
d_R	3	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	$-\frac{1}{3}$	$+$
L_L	3	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	0	$+$
e_R	3	$\mathbf{1}$	$\mathbf{1}$	-1	0	$+$
ν_R	3	$\mathbf{1}$	$\mathbf{1}$	0	1	$+$
D_L	3	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	$\frac{2}{3}$	$-$
D_R	3	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	$-\frac{1}{3}$	$-$
$\tilde{H}_L$	3	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	0	$-$
$\tilde{H}_R$	3	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	-1	$-$
N_L	3	$\mathbf{1}$	$\mathbf{1}$	0	-1	$-$

Table 20. Particle content of the `U1Leptophobic`. The scalar and fermion fields are shown in the top and bottom of the table respectively.

The particle content is summarized in Table 20, while the scalar potential reads

$$\begin{aligned} V_{\text{scalar}} = {} & \tilde{m}_1^2|H_1|^2 + \tilde{m}_2^2|H_2|^2 + \frac{\lambda_1}{2}|H_1|^4 + \frac{\lambda_2}{2}|H_2|^4 + \lambda_3|H_1|^2|H_2|^2 + \lambda_4 H_1^\dagger H_2 H_2^\dagger H_1 \\ & + \tilde{m}_\Phi^2|\Phi|^2 + \frac{\lambda_\Phi}{2}|\Phi|^4 + \left(\mu_\Phi H_1^\dagger H_2 \Phi + \text{h.c.}\right) + \tilde{\lambda}_1|H_1|^2|\Phi|^2 + \tilde{\lambda}_2|H_2|^2|\Phi|^2\,, \end{aligned} \tag{2.21}$$

and the Yukawa interactions are given by

$$\begin{aligned} \mathcal{L}_Y = {} & y^u \overline{u_R} H_1^\dagger Q + y^d \overline{d_R} H_2 Q + y^e \overline{e_R} H_2 L + y^n \overline{n_R} H_1^\dagger L \\ & + y^D \overline{D_R} \Phi D_L + y^H_{ij} \overline{\widetilde{H}^j_R} \Phi \widetilde{H}^i_L + y^N_{ij} \overline{N^{c\,i}_L} H_1^\dagger \widetilde{H}^j_L + y'^N_{ij} \overline{\widetilde{H}^i_R} H_2 N^j_L + \text{h.c.}\,. \end{aligned} \tag{2.22}$$

We assume the following symmetry breaking pattern

$$H_{1/2} = \frac{1}{\sqrt{2}} \begin{pmatrix} v_{1/2} + \phi_{1/2} + i\sigma_{1/2} \\ 0 \end{pmatrix}\,, \qquad \Phi = \frac{1}{\sqrt{2}}\left(v_\Phi + \phi_\Phi + i\sigma_\Phi\right)\,, \tag{2.23}$$

where once again ϕ_i and σ_i are the CP-even and CP-odd components, respectively, and we define $\tan\beta \equiv \frac{v_2}{v_1}$.

The 750 GeV candidate is taken to be dominantly composed by ϕ_Φ, that is, the real CP-even degree of freedom of the Φ field after $U(1)_X$ symmetry breaking. It couples to photons and gluons via loops of the new fermions. The model cannot explain the diphoton excess with Yukawa couplings in the perturbative range, but the authors use values between 5 and 10. As stressed in [1], this renders the perturbative calculation, and hence the results, very questionable.

2.2.5 $U(1)'$ extension with a Z' mimicking a scalar resonance

- **Reference:** [25]
- **Model name:** `U1Extensions/trickingLY`

The idea of this model is that the extra neutral gauge boson decays into $S\gamma$, whereas the scalar S itself decays into a diphoton final state. Because of the high boost, the two photons from S appear to be a single photon in the detector.

Ref. [25] works in a model realization where the third-generation quarks are charged under $U(1)'$ whereas the first and second generations are not. While that can be viewed as a toy model to make a point, an actual realisation would either need additional Higgs representations or flavour-universal $U(1)'$ charges in order to reproduce the correct CKM matrix. Consequently, we will work with $U(1)'$ charges for all three generations of SM quarks and also have to use three generations of each additional exotic particle for anomaly cancellation.

The particle content of the model is summarized in Table 21. For anomaly cancellation, the condition $Q_1 - Q_2 = -3$ must hold. Furthermore, there are four different valid choices for the hypercharge assignments Y_i [25]:

$$(Y_1,\, Y_2,\, Y_3) = (a, a + \frac{1}{2}, a - \frac{1}{2})\,. \tag{2.24}$$

In the model implementation at hand, we choose $(Y_1,\, Y_2,\, Y_3) = (\frac{1}{2}, 1, 0)$.

The Yukawa interactions including fields beyond the SM read

$$\mathcal{L} \supset Y_v\, \bar{\nu}_R\, \tilde{H} L_L + \eta_1\, \bar{X}_R\, S^*\, X_L + \eta_2 \bar{R}_R\, S\, R_L + \eta_3\, \bar{\xi}_R\, S\, \xi_L\,. \tag{2.25}$$

The scalar potential is given by

$$V = -\mu^2|H|^2 - \mu_S^2|S|^2 + \lambda_H|H|^4 + \lambda_S|S|^4 + \lambda_{HS}|S|^2|H|^2\,, \tag{2.26}$$

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)'$
H	1	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	0
S	1	$\mathbf{1}$	$\mathbf{1}$	0	$Q_1 - Q_2 = -3$
Q_L	3	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$	1
u_R	3	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	1
d_R	3	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	1
L_L	3	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	0
l_R	3	$\mathbf{1}$	$\mathbf{1}$	-1	0
ν_R	3	$\mathbf{1}$	$\mathbf{1}$	0	0
X_L	3	$\mathbf{3}$	$\mathbf{2}$	Y_1	Q_1
X_R	3	$\mathbf{3}$	$\mathbf{2}$	Y_1	Q_2
R_L	3	$\mathbf{1}$	$\mathbf{1}$	Y_2	Q_2
R_R	3	$\mathbf{1}$	$\mathbf{1}$	Y_2	Q_1
ξ_L	3	$\mathbf{1}$	$\mathbf{1}$	Y_3	Q_2
ξ_R	3	$\mathbf{1}$	$\mathbf{1}$	Y_3	Q_1

Table 21. Fermionic and scalar particle content of the `trickingLY`. Here $X_L = (x_{1L}, x_{2L})$, $X_R = (x_{1R}, x_{2R})$ and $H = (H^0, H^-)$.

where the $U(1)'$ symmetry is broken spontaneously as soon as S receives a VEV according to $\langle S \rangle = v_S/\sqrt{2}$.

Unfortunately, the most interesting vertex for this model, $Z' - S - \gamma$, only arises at the loop level. This would require a handling via an effective operator which is currently not supported in the automatized tools advertised. Hence, it is not possible to recast the results of Ref. [25] based on this model implementation only.

Note that, in principle, the decay $Z' \to ZS$ is already possible at tree level due to $Z - Z'$ mixing and dominates over the decay into $S\gamma$. Therefore, in order to achieve the desired effect, the mixed gauge interaction term $F^{\mu\nu}F'_{\mu\nu}$ must be forbidden while allowing for $SF^{\mu\nu}F'_{\mu\nu}$ which is hard to justify in general.

3 Non-abelian gauge-group extensions of the SM

3.1 Left-right symmetric models

Left-right (LR) symmetric models can potentially provide an interesting explanation of the diphoton excess through the use of an extended scalar sector that is necessary to spontaneously break the enlarged gauge group $\mathcal{G}_{L-R} \equiv SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ to the SM gauge group. However, due to the large number of fields, these models are often difficult to analyse even at tree-level. Therefore as a starting point we provide model files for four different left-right models that have been proposed in the literature to explain the diphoton excess. These four models are based on the above mentioned gauge group $\mathcal{G}_{L-R}$ with, in two cases, further Abelian gauged or global symmetries.

3.1.1 Left-right symmetric model without bi-doublets

- **Reference:** [26–28] (see also Ref. [62])
- **Model name:** `LRmodels/LR-VL`

In Ref. [26] the authors explored the possibility of explaining the observed diphoton excess in the context of the minimal left-right symmetric model. They show that it is not possible and that an alternative model is necessary. Therefore they give up on standard Yukawa couplings, for which bi-doublets are necessary, and consider separate $SU(2)_L$- and $SU(2)_R$-Higgs-doublets. In order to be able to introduce Yukawa interactions, new vectorlike $SU(2)_i$-singlet fermions are introduced. After integrating out the vectorlike fermions, the SM fermion masses are generated through a universal seesaw mechanism [63, 64].

The particle content for the model is shown in Table 22. Note that the authors consider the second-lightest CP-even Higgs, which should be predominantly singlet-like, as the particle responsible for the observed resonance. The scalar potential given the particle content and consistent with the symmetries $(SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L})$ is

$$\begin{aligned} V = {} & M_S^2 S^2 + \left(\mu_L^2 - \mu_{SL} S\right)|H_L|^2 + \left(\mu_R^2 + \mu_{SR} S\right)|H_R|^2 - \lambda_S S^4 - \lambda_L |H_L|^4 - \lambda_R |H_R|^4 \\ & - \lambda_{LR}|H_L|^2|H_R|^2 - \lambda_{SL} S^2 |H_L|^2 - \lambda_{SR} S^2 |H_R|^2 \,. \end{aligned} \tag{3.1}$$

The Yukawa interactions can be written as[4]

$$\begin{aligned} \mathcal{L}_Y = {} & Y_U \left(\overline{q_L} H_L U_R + \overline{q_R} H_R U_L\right) + Y_D \left(\overline{q_L} \tilde{H}_L D_R + \overline{q_R} \tilde{H}_R D_L\right) + Y_E \left(\overline{l_L} \tilde{H}_L E_R + \overline{l_R} \tilde{H}_R E_L\right) \\ & + Y_N \left(\overline{l_L} H_L N_R + \overline{l_R} H_R N_L\right) + \frac{1}{2} m_{NM} \left(\overline{N_R^c} N_R + \overline{N_L^c} N_L\right) + m_{ND} \overline{N_L} N_R \\ & + \left(m_U + \lambda_U S\right) \overline{U_L} U_R + \left(m_D + \lambda_D S\right) \overline{D_L} D_R + \left(m_E + \lambda_E S\right) \overline{E_L} E_R + \text{h.c.}\,, \end{aligned} \tag{3.2}$$

[4] Note that our Yukawa interactions differ from the literature: in Ref. [26], they are defined as, e.g., $\overline{q_L} H_L^\dagger U_L$ which contracts to zero because of the implicit left/right projection operators. Moreover, in Refs. [26, 27] the 'conjugate' assignments of the $H_{L/R}$ need to be exchanged in order to obtain a gauge-invariant Lagrangian.

Field	Generations	$SU(3)_C$	$SU(2)_L$	$SU(2)_R$	$U(1)_{B-L}$
H_L	1	$\mathbf{1}$	$\mathbf{2}$	$\mathbf{1}$	$-\frac{1}{2}$
H_R	1	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$
S	1	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	0
$q_L = (u_{L,i}, d_{L,i})^T$	3	$\mathbf{3}$	$\mathbf{2}$	$\mathbf{1}$	$\frac{1}{6}$
$q_R = (u_{R,i}, d_{R,i})^T$	3	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{6}$
$l_L = (\nu_{L,i}, e_{L,i})^T$	3	$\mathbf{1}$	$\mathbf{2}$	$\mathbf{1}$	$-\frac{1}{2}$
$l_R = (\nu_{R,i}, e_{R,i})^T$	3	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$
U_L	3	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{1}$	$\frac{2}{3}$
U_R	3	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{1}$	$\frac{2}{3}$
D_L	3	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{1}$	$-\frac{1}{3}$
D_R	3	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{1}$	$-\frac{1}{3}$
$E_{L/R}$	3	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	-1
$N_{L/R}$	3	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	0

Table 22. Matter content and charge assignments for the `LR-VL` model. The scalar/fermionic fields are shown in the top/bottom of the table respectively. The generation index i runs over $i = 1, 2, 3$.

where N_L^c is the charge conjugate of N_L. Note that we have included both Majorana and Dirac mass terms for the fermionic singlet $N_{L/R}$. We assume the following symmetry breaking VEVs

$$\langle H_L \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v_L \\ 0 \end{pmatrix} , \qquad \langle H_R \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v_R \\ 0 \end{pmatrix} , \qquad \langle S \rangle = v_S \, . \tag{3.3}$$

3.1.2 Left-right symmetric model with $U(1)_L \times U(1)_R$

- **Reference:** [29]
- **Model name:** `LRmodels/LRLR`

This model is based on the gauge group $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_L \times U(1)_R$. The inclusion of extra vectorlike $SU(2)$-singlet fermions allows for the generation of the SM fermion masses via a see-saw mechanism. Additionally, a parity symmetry is imposed to ensure a vanishing $\bar{\theta}$ parameter at tree-level in the QCD Lagrangian [65], in order to solve the strong CP-problem without introducing an axion.

The particle content of this model and charge assignments under the gauge symmetries are shown in Table 23. The proposed candidate for the 750 GeV resonance is taken to be one of the CP-even scalars associated with the $SU(2)$-singlet Higgs scalars σ_D, σ_U and σ_E that are responsible for the breaking $U(1)_L \times U(1)_R \to U(1)_{B-L}$. The decays of this state into digluon and diphoton final states are assumed to proceed via loops involving the $SU(2)$-singlet fermions.

Field	$SU(3)_c$	$SU(2)_L$	$SU(2)_R$	$U(1)_L$	$U(1)_R$
σ_U	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$\frac{2}{3}$	$-\frac{2}{3}$
σ_D	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$-\frac{1}{3}$	$\frac{1}{3}$
σ_E	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	-1	1
ϕ_L	$\mathbf{1}$	$\mathbf{2}$	$\mathbf{1}$	$-\frac{1}{2}$	0
ϕ_R	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{2}$	0	$-\frac{1}{2}$
Δ_L	$\mathbf{1}$	$\mathbf{3}$	$\mathbf{1}$	1	0
Δ_R	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{3}$	0	1
$q_{L,i} = (u_{L,i}, d_{L,i})^T$	$\mathbf{3}$	$\mathbf{2}$	$\mathbf{1}$	$\frac{1}{6}$	0
$q_{R,i} = (u_{R,i}, d_{R,i})^T$	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{2}$	0	$\frac{1}{6}$
$l_{L,i} = (\nu_{L,i}, e_{L,i})^T$	$\mathbf{1}$	$\mathbf{2}$	$\mathbf{1}$	$-\frac{1}{2}$	0
$l_{R,i} = (\nu_{R,i}, e_{R,i})^T$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{2}$	0	$-\frac{1}{2}$
$U_{L,i}$	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{1}$	$-\frac{2}{3}$	0
$U_{R,i}$	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{1}$	0	$-\frac{2}{3}$
$D_{L,i}$	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{1}$	$\frac{1}{3}$	0
$D_{R,i}$	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{1}$	0	$\frac{1}{3}$
E_L	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	1	0
E_R	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	0	1

Table 23. The $SU(3)_c \times SU(2)_L \times SU(2)_R \times U(1)_L \times U(1)_R$ charge assignments for the scalar/fermions in the `LRLR` shown in the top/bottom of the table. The generation index i runs over $i = 1, 2, 3$.

The Yukawa interactions consistent with the imposed parity are given by

$$\begin{aligned} -\mathcal{L} = {} & y_U \left(\bar{q}_L \phi_L U_L^c + \bar{q}_R \phi_R U_R^c\right) + f_U \sigma_U^* \bar{U}_L U_R + y_D \left(\bar{q}_L \tilde{\phi}_L D_L^c + \bar{q}_R \tilde{\phi}_R D_R^c\right) + f_D \sigma_D^* \bar{D}_L D_R \\ & + y_E \left(\bar{l}_L \tilde{\phi}_L E_L^c + \bar{l}_R \tilde{\phi}_R E_R^c\right) + f_E \sigma_E^* \bar{E}_L E_R + Y_L \left(\bar{l}_L^c i\tau_2 \Delta_L l_L + \bar{l}_R^c i\tau_2 \Delta_R l_R\right) + \text{h.c.} \end{aligned} \tag{3.4}$$

The parity symmetry is taken to be softly broken, so that the part of the scalar potential considered in Ref. [29] is given by

$$
\begin{aligned}
V = &\lambda \left(\sigma_E^* \sigma_D^3 + \text{h.c.}\right) + \xi \left(\sigma_U \sigma_D^2 + \text{h.c.}\right) \\
&+ \eta \left(\sigma_E \sigma_D^* \sigma_U + \text{h.c.}\right) + \mu_{\phi_L}^2 \phi_L^\dagger \phi_L + \mu_{\phi_R}^2 \phi_R^\dagger \phi_R \\
&+ \mu_{\Delta_L}^2 \text{Tr}\left(\Delta_L^\dagger \Delta_L\right) + \mu_{\Delta_R}^2 \text{Tr}\left(\Delta_R^\dagger \Delta_R\right) \\
&+ \rho_L \left(\phi_L^T i\tau_2 \Delta_L \phi_L + \text{h.c.}\right) + \rho_R \left(\phi_R^T i\tau_2 \Delta_R \phi_R + \text{h.c.}\right) .
\end{aligned} \tag{3.5}
$$

The couplings and masses λ, ξ, η, $\mu_{\phi_L}^2$, $\mu_{\phi_R}^2$, $\mu_{\Delta_L}^2$, $\mu_{\Delta_R}^2$, ρ_L and ρ_R are taken to be real, with $\mu_{\phi_L}^2 \neq \mu_{\phi_R}^2$, $\mu_{\Delta_L}^2 \neq \mu_{\Delta_R}^2$ and $\rho_L \neq \rho_R$. Note that Eq. (3.4) and Eq. (3.5) differ from Eq. (6) and Eq. (7) in Ref. [29], which as given are not gauge invariant. One may also include a large number of additional terms that are allowed by the gauge symmetries, given by

$$
\begin{aligned}
V' = &\kappa \left(\sigma_D \sigma_E \sigma_U^2 + \text{h.c.}\right) + \mu_D^2 |\sigma_D|^2 + \lambda_{DD} |\sigma_D|^4 + \mu_U^2 |\sigma_U|^2 + \lambda_{UU} |\sigma_U|^4 + \mu_E^2 |\sigma_E|^2 \\
&+ \lambda_{EE} |\sigma_E|^4 + \lambda_{DU} |\sigma_D|^2 |\sigma_U|^2 + \lambda_{DE} |\sigma_D|^2 |\sigma_E|^2 + \lambda_{UE} |\sigma_U|^2 |\sigma_E|^2 + \lambda_{LL} |\phi_L|^4 \\
&+ \lambda_{RR} |\phi_R|^4 + \lambda_{LR} |\phi_L|^2 |\phi_R|^2 + \rho_{R_1} \text{Tr}(\Delta_R^\dagger \Delta_R) \text{Tr}(\Delta_R^\dagger \Delta_R) + \rho_{R_2} \text{Tr}(\Delta_R \Delta_R) \text{Tr}(\Delta_R^\dagger \Delta_R^\dagger) \\
&+ \rho_{L_1} \text{Tr}(\Delta_L^\dagger \Delta_L) \text{Tr}(\Delta_L^\dagger \Delta_L) + \rho_{L_2} \text{Tr}(\Delta_L \Delta_L) \text{Tr}(\Delta_L^\dagger \Delta_L^\dagger) + \rho_3 \text{Tr}(\Delta_L^\dagger \Delta_L) \text{Tr}(\Delta_R^\dagger \Delta_R) \\
&+ \eta_{LL} |\phi_L|^2 \text{Tr}(\Delta_L^\dagger \Delta_L) + \eta_{RL} |\phi_R|^2 \text{Tr}(\Delta_L^\dagger \Delta_L) + \eta_{LR} |\phi_L|^2 \text{Tr}(\Delta_R^\dagger \Delta_R) + \eta_{RR} |\phi_R|^2 \text{Tr}(\Delta_R^\dagger \Delta_R) \\
&+ e_{RR_1} \phi_R^\dagger \Delta_R^\dagger \Delta_R \phi_R - e_{RR_2} \phi_R^\dagger \Delta_R \Delta_R^\dagger \phi_R + e_{LL_1} \phi_L^\dagger \Delta_L^\dagger \Delta_L \phi_L - e_{LL_2} \phi_L^\dagger \Delta_L \Delta_L^\dagger \phi_L \\
&+ \sum_{f=U,D,E} |\sigma_f|^2 \left[\lambda_{fL} |\phi_L|^2 + \lambda_{fR} |\phi_R|^2 + \tilde{\lambda}_{fL} \text{Tr}\left(\Delta_L^\dagger \Delta_L\right) + \lambda_{fR} \text{Tr}\left(\Delta_R^\dagger \Delta_R\right)\right] .
\end{aligned} \tag{3.6}
$$

The full scalar potential that we consider is then $V + V'$.

The $SU(2)$-singlet Higgs scalars are assumed to acquire VEVs of the form

$$
\langle \sigma_D \rangle = \frac{v_D}{\sqrt{2}} , \qquad \langle \sigma_U \rangle = \frac{v_U}{\sqrt{2}} , \qquad \langle \sigma_E \rangle = \frac{v_E}{\sqrt{2}} . \tag{3.7}
$$

resulting in the breaking $U(1)_L \times U(1)_R \to U(1)_{B-L}$. The $SU(2)$-doublet Higgs scalars,

$$
\phi_{L,R} = \begin{pmatrix} \phi_{L,R}^0 \\ \phi_{L,R}^- \end{pmatrix} , \tag{3.8}
$$

are taken to develop VEVs of the form

$$
\langle \phi_{L,R} \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} v_{L,R} \\ 0 \end{pmatrix} . \tag{3.9}
$$

The non-zero VEV of ϕ_R^0 leads to the breakdown of $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ to $SU(2)_L \times U(1)_Y$, which is subsequently broken by the VEV of ϕ_L^0. As a result, the triplet scalars

$$\Delta_{L,R} = \begin{pmatrix} \frac{\delta^+_{L,R}}{\sqrt{2}} & \delta^{++}_{L,R} \\ \delta^0_{L,R} & -\frac{\delta^+_{L,R}}{\sqrt{2}} \end{pmatrix} \tag{3.10}$$

also acquire VEVs of the form

$$\langle \Delta_{L,R} \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & 0 \\ u_{L,R} & 0 \end{pmatrix} . \tag{3.11}$$

3.1.3 Left-right symmetric model with fermionic and scalar triplets

- **Reference:** [30]
- **Model name:** `LRmodels/tripletLR`

In this model the diphoton signal is produced through a cascade decay, namely, $pp \to Z' \to XY \to XX(\delta^0 \to \gamma\gamma)$ where X and Y are unspecified soft states and δ_0 is the neutral component of the $SU(2)_R$ scalar triplet. However, in order to sufficiently boost the rate three $SU(2)_R$ triplet fermion fields are added to the model, T_1, T_2 and T_3. The model is based on a $SU(3)_C \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge group which is broken to the SM gauge group through the VEV of the triplet field Δ_R whereby EWSB proceeds through the bi doublet Φ VEVs. The entire particle content of the model is illustrated in Table 24.

The scalar fields of the model can be expressed as

$$\Phi = \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix} , \quad \text{and} \quad \Delta_R = \begin{pmatrix} \frac{\delta_R}{\sqrt{2}} & \delta_R^{++} \\ \delta_R^0 & -\frac{\delta_R^+}{\sqrt{2}} \end{pmatrix} , \tag{3.12}$$

leading to a scalar potential of the form

$$\begin{aligned} V = {} & \mu_1^2 \mathrm{Tr}\left(\Phi^\dagger \Phi\right) + \mu_2^2 \left[\mathrm{Tr}(\widetilde{\Phi}\Phi^\dagger) + \mathrm{Tr}(\widetilde{\Phi}^\dagger \Phi)\right] + \mu_3^2 \mathrm{Tr}\left(\Delta_R \Delta_R^\dagger\right) + \lambda_1 \left[\mathrm{Tr}\left(\Phi^\dagger \Phi\right)\right]^2 \\ & + \lambda_2 \left\{\left[\mathrm{Tr}(\widetilde{\Phi}\Phi^\dagger)\right]^2 + \left[\mathrm{Tr}(\widetilde{\Phi}^\dagger \Phi)\right]^2\right\} + \lambda_3 \mathrm{Tr}(\widetilde{\Phi}\Phi^\dagger)\mathrm{Tr}(\widetilde{\Phi}^\dagger \Phi) \\ & + \lambda_4 \mathrm{Tr}(\Phi^\dagger \Phi)\left[\mathrm{Tr}(\widetilde{\Phi}\Phi^\dagger) + \mathrm{Tr}(\widetilde{\Phi}^\dagger \Phi)\right] + \rho_1 \left[\mathrm{Tr}(\Delta_R \Delta_R^\dagger)\right]^2 + \rho_2 \mathrm{Tr}(\Delta_R \Delta_R)\mathrm{Tr}(\Delta_R^\dagger \Delta_R^\dagger) \\ & + \alpha_1 \mathrm{Tr}\left(\Phi^\dagger \Phi\right)\mathrm{Tr}(\Delta_R \Delta_R^\dagger) + \left\{\alpha_2 e^{i\delta} \mathrm{Tr}(\widetilde{\Phi}^\dagger \Phi)\mathrm{Tr}(\Delta_R \Delta_R^\dagger) + \text{h.c.}\right\} + \alpha_3 \mathrm{Tr}(\Phi \Phi^\dagger \Delta_R \Delta_R^\dagger) , \end{aligned} \tag{3.13}$$

Field	Generations	$SU(3)_C$	$SU(2)_L$	$SU(2)_R$	$U(1)_{B-L}$
Φ	1	$\mathbf{1}$	$\mathbf{2}$	$\mathbf{2}$	0
Δ_R	1	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{3}$	1
$Q_L = (u_{L,i}, d_{L,i})^T$	3	$\mathbf{3}$	$\mathbf{2}$	$\mathbf{1}$	$\frac{1}{6}$
$Q_R = (u_{R,i}, d_{R,i})^T$	3	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{6}$
$L_L = (\nu_{L,i}, e_{L,i})^T$	3	$\mathbf{1}$	$\mathbf{2}$	$\mathbf{1}$	$-\frac{1}{2}$
$L_R = (\nu_{R,i}, e_{R,i})^T$	3	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$
T_1	1	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{3}$	0
T_2	1	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{3}$	1
T_3	1	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{3}$	-1

Table 24. Matter content and charge assignments for the `tripletLR` model. The scalar/fermionic fields are shown in the top/bottom of the table respectively. The generation index i runs over $i = 1, 2, 3$.

where $\tilde{\Phi} \equiv -\sigma_2 \Phi^* \sigma_2$. The Yukawa interactions of the model are given by

$$\begin{aligned}\mathcal{L}_Y &= Y_1^\alpha \overline{\Psi}_L \Phi \Psi_R + Y_2^\alpha \overline{\Psi}_L \widetilde{\Phi} \Psi_R + Y_{DR} L_R^T \mathcal{C} (i\sigma_2) \Delta_R L_R \\ &+ \frac{1}{2} m_1 \mathrm{Tr}\,(T_1 T_1) + m_{23} \mathrm{Tr}\,(T_2 T_3) + \lambda_{T_{13}} \mathrm{Tr}\,(T_1 T_3 \Delta_R) + \lambda_{T_{12}} \mathrm{Tr}(T_1 T_2 \Delta_R^\dagger)\,, \end{aligned} \tag{3.14}$$

where α runs over the quarks and leptons $\alpha = Q, L$ and $\Psi_{L,R} = (\psi^u_{L,R}, \psi^d_{L,R})$ with $\psi^u = u, \nu$ and $\psi^d = d, \ell$. The triplet fermion fields can be written as bidoublets

$$T_1 = \begin{pmatrix} \frac{t_1^0}{\sqrt{2}} & t_1^+ \\ t_1^- & -\frac{t_1^0}{\sqrt{2}} \end{pmatrix}, \qquad T_2 = \begin{pmatrix} \frac{t_2^+}{\sqrt{2}} & t_2^{++} \\ t_2^0 & -\frac{t_2^+}{\sqrt{2}} \end{pmatrix}, \qquad T_3 = \begin{pmatrix} \frac{t_3^-}{\sqrt{2}} & t_3^0 \\ t_3^{--} & -\frac{t_3^-}{\sqrt{2}} \end{pmatrix}. \tag{3.15}$$

Finally, $\mathcal{C}$ is the charge conjugation operator. The VEVs of the scalar fields in the model take the form

$$\langle \Phi \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} \kappa_1 & 0 \\ 0 & \kappa_2 \end{pmatrix} \quad \text{and} \quad \langle \Delta_R \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 & 0 \\ v_R & 0 \end{pmatrix}. \tag{3.16}$$

3.1.4 Dark left-right symmetric model

- **Reference:** [31]
- **Model name:** `LRmodels/darkLR`

The main idea of this model is to add an additional symmetry in order to stabilize the DM candidate, namely right-handed neutrinos, so that they cannot decay via a

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$SU(2)_R$	$U(1)_{B-L}$	$U(1)_S$	Lep.
Φ	1	$\mathbf{1}$	$\mathbf{2}$	$\mathbf{2}$	0	$\frac{1}{2}$	0
Δ_L	1	$\mathbf{1}$	$\mathbf{3}$	$\mathbf{1}$	1	-2	0
Δ_R	1	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{3}$	1	-1	0
H_L	1	$\mathbf{1}$	$\mathbf{2}$	$\mathbf{1}$	$\frac{1}{2}$	0	0
H_R	1	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{2}$	$-\frac{1}{2}$	0
$Q_L = (u_{L,i}, d_{L,i})^T$	3	$\mathbf{3}$	$\mathbf{2}$	$\mathbf{1}$	$\frac{1}{6}$	0	$(0,0)$
$Q_R = (u_{R,i}, d_{R,i})^T$	3	$\mathbf{3}$	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{6}$	$\frac{1}{2}$	$(0,1)$
$L_L = (\nu_{L,i}, e_{L,i})^T$	3	$\mathbf{1}$	$\mathbf{2}$	$\mathbf{1}$	$-\frac{1}{2}$	1	$(1,1)$
$L_R = (\nu_{R,i}, e_{R,i})^T$	3	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	$\frac{1}{2}$	$(0,1)$
d_R	3	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$-\frac{1}{3}$	0	0
x_L	3	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$	$-\frac{1}{3}$	1	1

Table 25. Matter content and charge assignments for the `darkLR` model. The scalar/fermionic fields are shown in the top/bottom of the table respectively. The generation index, denoted Gen. in the table, i runs over $i = 1, 2, 3$. Additionally the model includes a lepton number symmetry where the quantum numbers are denoted with Lep. above.

W' channel. This additional symmetry takes the form of a global Abelian symmetry labelled as $U(1)_S$. The spontaneous breaking of $SU(2)_R \times U(1)_S$ is such that the combination $\tilde{L} = S - T_{3R}$, where T_{3R} is the third component of the right-handed isospin, remains unbroken. Here $\tilde{L}$ is interpreted as a generalised lepton number.

The full particle content of the model is shown in Table 25. Note that the scalar sector of the model is enlarged to include both left- and right-handed triplets and doublets as well as the usual bi-doublet. These scalar fields can be expressed as

$$\Phi = \begin{pmatrix} \phi_1^0 & \phi_2^+ \\ \phi_1^- & \phi_2^0 \end{pmatrix} , \quad \Phi_{L,R} = \begin{pmatrix} H_{L,R}^+ \\ \phi_{L,R}^0 \end{pmatrix} , \quad \Delta_{L,R} = \begin{pmatrix} \frac{\delta_{L,R}}{\sqrt{2}} & \delta_{L,R}^{++} \\ \delta_{L,R}^0 & -\frac{\delta_{L,R}^+}{\sqrt{2}} \end{pmatrix} . \tag{3.17}$$

Subsequently, the proposed candidate for the diphoton excess is ϕ_R^0, the neutral component of the $SU(2)_R$ doublet. Running in the loop will be W'-bosons, as well as δ_R^+ and δ_R^{++} Higgses. However, these particles are insufficient to boost the rate to diphotons, therefore additional quarks x_L and d_R are introduced.

The scalar potential of the model is

$$
\begin{aligned}
V &= \mu_1^2 \mathrm{Tr}\left(\Phi^\dagger\Phi\right) + \mu_{TR}^2 \mathrm{Tr}\left(\Delta_R\Delta_R^\dagger\right) + \mu_{TL}^2 \mathrm{Tr}\left(\Delta_L\Delta_L^\dagger\right) \\
&+ \mu_{DL}^2 H_L^\dagger H_L + \mu_{DR}^2 H_R^\dagger H_R + \lambda_1 \left[\mathrm{Tr}\left(\Phi^\dagger\Phi\right)\right]^2 + \lambda_3 \mathrm{Tr}(\widetilde{\Phi}\Phi^\dagger)\mathrm{Tr}(\widetilde{\Phi}^\dagger\Phi) \\
&+ \rho_1 \left\{\left[\mathrm{Tr}(\Delta_L\Delta_L^\dagger)\right]^2 + \left[\mathrm{Tr}(\Delta_R\Delta_R^\dagger)\right]^2\right\} + \beta_2\left[\mathrm{Tr}(\widetilde{\Phi}\Delta_R\Phi^\dagger\Delta_L^\dagger) + \mathrm{Tr}(\widetilde{\Phi}^\dagger\Delta_L\Phi\Delta_R^\dagger)\right] \\
&+ \rho_2\left\{\mathrm{Tr}(\Delta_L\Delta_L)\mathrm{Tr}(\Delta_L^\dagger\Delta_L^\dagger) + \mathrm{Tr}(\Delta_R\Delta_R)\mathrm{Tr}(\Delta_R^\dagger\Delta_R^\dagger)\right\} + \rho_3 \mathrm{Tr}(\Delta_L\Delta_L^\dagger)\mathrm{Tr}(\Delta_R\Delta_R^\dagger) \\
&+ \alpha_1 \mathrm{Tr}\left(\Phi^\dagger\Phi\right)\left[\mathrm{Tr}(\Delta_L\Delta_L^\dagger) + \mathrm{Tr}(\Delta_R\Delta_R^\dagger)\right] + \alpha_3\left[\mathrm{Tr}(\Phi\Phi^\dagger\Delta_L\Delta_L^\dagger) + \mathrm{Tr}(\Phi\Phi^\dagger\Delta_R\Delta_R^\dagger)\right] \\
&+ \eta_{LL} H_L^\dagger H_L \Delta_L\Delta_L^\dagger + \eta_{RL} H_R^\dagger H_R \Delta_L\Delta_L^\dagger + \eta_{LR} H_L^\dagger H_L \Delta_R\Delta_R^\dagger + \eta_{RR_1} H_R^\dagger H_R \Delta_R\Delta_R^\dagger \\
&+ \eta_{RR_2} H_R^\dagger \Delta_R^\dagger \Delta_R H_R + \eta_{RR_3} H_R^\dagger \Delta_R \Delta_R^\dagger H_R + \lambda_L |H_L|^4 + \lambda_R |H_R|^4 \lambda_{LR} |H_L|^2 |H_R|^2 \\
&+ \beta_L |H_L|^2 \Phi^\dagger\Phi + \beta_R |H_R|^2 \Phi^\dagger\Phi \\
&+ \left\{\alpha_4 H_L^\dagger \Phi \Delta_R H_R^\dagger + \xi_R \tilde{H}_R^\dagger \Delta_R^\dagger H_R + \xi_{LR} H_R \Phi H_L^\dagger + \mathrm{h.c.}\right\} ,
\end{aligned} \tag{3.18}
$$

where $\tilde{\Phi} \equiv -\sigma_2\Phi^*\sigma_2$. Note that there are a number of extra terms present in this potential compared to [31], which are allowed under the symmetries of the model. Finally the Yukawa interactions are given by

$$
\begin{aligned}
-\mathcal{L}_Y &= Y_{L_1}\overline{L}_L\Phi L_R + Y_{Q_1}\overline{Q}_L\tilde{\Phi}Q_R + Y_{Q_2}\overline{Q}_L H_L d_R + Y_{Q_3}\overline{x}_L\tilde{H}_R Q_R \\
&+ \frac{1}{2}Y_{DL}L_L^T\mathcal{C}(i\sigma_2)\Delta_L L_L + \frac{1}{2}Y_{DR}L_R^T\mathcal{C}(i\sigma_2)\Delta_R L_R + \mathrm{h.c.}\,,
\end{aligned} \tag{3.19}
$$

where $\mathcal{C}$ is the charge conjugation operator. The structure of the VEVs of the model are

$$
\langle\Phi\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix}\kappa_1 & 0 \\ 0 & \kappa_2\end{pmatrix} \qquad \langle\Phi_{L,R}\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix}0 \\ v_{L,R}^D\end{pmatrix}, \quad \langle\Delta_{L,R}\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix}0 & 0 \\ v_{L,R}^T & 0\end{pmatrix}. \tag{3.20}
$$

3.2 331 models

Models based on the $\mathrm{SU(3)_c} \times \mathrm{SU(3)_L} \times \mathrm{U(1)}_{\mathcal{X}}$ gauge symmetry [66–72], 331 for short, constitute an extension of the SM that could explain the number of generations of matter fields. This is possible as anomaly cancellation forces the number of generations to be equal to the number of quark colours.

Regarding the diphoton excess, 331 models automatically include all the required ingredients to explain the hint. First, the usual $SU(2)_L$ Higgs doublet must be promoted to a $SU(2)_L$ triplet, the new component being a singlet under the standard $\mathrm{SU(3)_c} \times \mathrm{SU(2)_L} \times \mathrm{U(1)_Y}$ symmetry. Similarly, the group structure requires the introduction of new coloured fermions to complete the $SU(3)_L$ quark multiplets, these exotic quarks being $\mathrm{SU}(3)_\mathrm{c} \times \mathrm{SU}(2)_\mathrm{L} \times \mathrm{U}(1)_\mathrm{Y}$ vector-like singlets after the breaking of $\mathrm{SU}(3)_\mathrm{c} \times \mathrm{SU}(3)_\mathrm{L} \times \mathrm{U}(1)_\mathcal{X}$. Therefore, $\mathrm{SU}(3)_\mathrm{c} \times \mathrm{SU}(3)_\mathrm{L} \times \mathrm{U}(1)_\mathcal{X}$ models naturally embed the simple *singlet + vector-like fermions* framework proposed to explain the diphoton excess.

There are several variants of $\mathrm{SU}(3)_\mathrm{c} \times \mathrm{SU}(3)_\mathrm{L} \times \mathrm{U}(1)_\mathcal{X}$ models. These are characterized by their β parameter [5], which defines the electric charge operator as [6]

$$Q = T_3 + \beta\, T_8 + \mathcal{X} \, . \tag{3.21}$$

First, in Section 3.2.1 we consider the model in Ref. [32]. This 331 variant has $\beta = 1/\sqrt{3}$, which fixes the electric charges of all the states contained in the $SU(2)_L$ triplets and anti-triplets to the usual $0, \pm 1$ values. In Section 3.2.2 we consider a 331 model with $\beta = -\sqrt{3}$, a value leading to exotic electric charges. This 331 variant has been discussed in the context of the diphoton excess in [33, 74, 75]. Although the mechanism to explain the diphoton excess is exactly the same as in [32], the presence of the exotic states leads to slightly different numerical results.

On the $SU(3)$ generators in `SARAH`

The most common choice for the $SU(3)$ generators is $T_a = \frac{\lambda_a}{2}$, with λ_a $(a = 1, \ldots, 8)$ the Gell-Mann matrices. However, this is just one of the possible representations. In fact, `SARAH` uses a different set of matrices, $T_a^{\texttt{SARAH}} = \frac{\Lambda_a}{2}$, following the conventions of `Susyno` [76]. The relation between the non-diagonal matrices in the two bases is

$$\lambda_1 = \Lambda_1 \, , \tag{3.22a}$$
$$\lambda_2 = \Lambda_4 \, , \tag{3.22b}$$
$$\lambda_4 = -\Lambda_6 \, , \tag{3.22c}$$
$$\lambda_5 = -\Lambda_3 \, , \tag{3.22d}$$
$$\lambda_6 = \Lambda_2 \, , \tag{3.22e}$$
$$\lambda_7 = \Lambda_5 \, . \tag{3.22f}$$

Concerning the diagonal matrices, the usual $\lambda_{3,8}$ Gell-Mann matrices,

$$\lambda_3 = \begin{pmatrix} 1 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & 0 \end{pmatrix} \quad , \quad \lambda_8 = \frac{1}{\sqrt{3}} \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -2 \end{pmatrix} \, , \tag{3.23}$$

[5] See [73] for a complete discussion of 331 models with generic β.

[6] Eq. (3.21) assumes that the $SU(3)$ generators are $T_a = \frac{\lambda_a}{2}$, with λ_a $(a = 1, \ldots, 8)$ the Gell-Mann matrices. However, this is not the convention used in `SARAH`, see below.

are replaced by $\Lambda_{7,8}$,

$$\Lambda_7 = \frac{1}{\sqrt{3}} \begin{pmatrix} 2 & 0 & 0 \\ 0 & -1 & 0 \\ 0 & 0 & -1 \end{pmatrix} \quad , \quad \Lambda_8 = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 1 & 0 \\ 0 & 0 & -1 \end{pmatrix} . \tag{3.24}$$

The electric charge operator can be written, using the conventions in `SARAH`, as

$$Q^{\texttt{SARAH}} = -T_8 - \beta\, T_7 + \mathcal{X} . \tag{3.25}$$

This in turn implies that the charge assignments in the $SU(3)$ multiplets must be adapted as well. For example, one can easily check that the electric charges of the first and third components of a $SU(3)$ triplet t are exchanged when going from the usual Gell-Mann representation to the basis choice employed in `SARAH`,

$$t = \begin{pmatrix} t_1 \\ t_2 \\ t_3 \end{pmatrix} \quad \longrightarrow \quad t^{\texttt{SARAH}} = \begin{pmatrix} t_3 \\ t_2 \\ t_1 \end{pmatrix} . \tag{3.26}$$

In the following we will use the standard conventions based on the Gell-Mann matrices in order to keep the discussion as close to the original works as possible. However, we emphasize that the implementation of the 331 models in `SARAH` requires this dictionary between the bases. It should also be noted that in the current implementation in `SARAH` of the 331 models described below, vertices involving four vector bosons in the generated model files for `CalcHep` cannot yet be handled correctly. In order to generate model files that will work with `CalcHep`, one must therefore exclude these vertices from being written out by `SARAH`.

3.2.1 331 model without exotic charges

- **Reference:** [32]
- **Model name:** `331/v1`

The model is based on the $\mathrm{SU}(3)_\mathrm{c} \times \mathrm{SU}(3)_\mathrm{L} \times \mathrm{U}(1)_\mathcal{X}$ gauge symmetry, extended with a global $U(1)_\mathcal{L}$ and an auxiliary $\mathbb{Z}_2$ symmetry to forbid some undesired couplings. The fermionic and scalar particle content of the model is summarized in Table 26. In addition, due to the extended group structure, the model contains 17 gauge bosons: the usual 8 gluons; 8 W_i bosons associated to $SU(3)_L$ and the B boson associated to $U(1)_\mathcal{X}$.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_{\mathcal{X}}$	$U(1)_{\mathcal{L}}$	$\mathbb{Z}_2$
Φ_1	1	$\mathbf{1}$	$\bar{\mathbf{3}}$	$\frac{2}{3}$	$\frac{2}{3}$	$+$
Φ_2	1	$\mathbf{1}$	$\bar{\mathbf{3}}$	$-\frac{1}{3}$	$-\frac{4}{3}$	$+$
Φ_3	1	$\mathbf{1}$	$\bar{\mathbf{3}}$	$-\frac{1}{3}$	$\frac{2}{3}$	$-$
Φ_X	1	$\mathbf{1}$	$\bar{\mathbf{3}}$	$-\frac{1}{3}$	$-\frac{4}{3}$	$+$
ψ_L	3	$\mathbf{1}$	$\bar{\mathbf{3}}$	$-\frac{1}{3}$	$-\frac{1}{3}$	$+$
e_R	3	$\mathbf{1}$	$\mathbf{1}$	-1	-1	$+$
s	3	$\mathbf{1}$	$\mathbf{1}$	0	1	$+$
$Q_L^{1,2}$	2	$\mathbf{3}$	$\mathbf{3}$	0	$-\frac{2}{3}$	$+$
Q_L^3	1	$\mathbf{3}$	$\bar{\mathbf{3}}$	$\frac{1}{3}$	$\frac{2}{3}$	$-$
u_R	3	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	0	$+$
T_R	1	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	0	$-$
d_R	3	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	0	$-$
D_R, S_R	2	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	0	$+$

Table 26. Fermionic and scalar particle content of the `331-v1` model. The scalar and fermion fields are shown in the top and bottom of the table respectively.

The fermionic $SU(3)_L$ triplets of the model can be decomposed as

$$\psi_L = \begin{pmatrix} \ell^- \\ -\nu \\ N^c \end{pmatrix}_L^{e,\mu,\tau}, \; Q_L^1 = \begin{pmatrix} u \\ d \\ D \end{pmatrix}_L, \; Q_L^2 = \begin{pmatrix} c \\ s \\ S \end{pmatrix}_L, \; Q_L^3 = \begin{pmatrix} b \\ -t \\ T \end{pmatrix}_L. \tag{3.27}$$

The notation used for the extra quarks that constitute the third components of the $SU(3)_L$ triplets $Q_L^{1,2,3}$ is motivated by the fact that their electric charges are $-1/3$ and $2/3$ for D/S and T, respectively. The scalar multiplets can be written as

$$\Phi_1 = \begin{pmatrix} \phi_1 \\ -\phi_1^- \\ S_1^- \end{pmatrix}, \quad \Phi_2 = \begin{pmatrix} \phi_2^+ \\ -\phi_2 \\ S_2 \end{pmatrix}, \quad \Phi_3 = \begin{pmatrix} \phi_3^+ \\ -\phi_3 \\ S_3 \end{pmatrix}, \quad \Phi_X = \begin{pmatrix} \phi_X^+ \\ -\phi_X \\ X \end{pmatrix}. \tag{3.28}$$

While ϕ_1^-, $\phi_{2,3}^+$ and S_1^- are electrically charged scalars, the components $\phi_{1,2,3,X}$, $S_{2,3}$ and X are neutral.

The Yukawa Lagrangian of the model can be split as

$$\mathcal{L}_Y = \mathcal{L}_Y^q + \mathcal{L}_Y^\ell\,, \tag{3.29}$$

where

$$\mathcal{L}_Y^q = \bar{Q}_L^{1,2}\, y^u u_R \Phi_1^* + \bar{Q}_L^3\, \tilde{y}^d d_R \Phi_1 + \bar{Q}_L^{1,2}\, \bar{y}^d \hat{d}_R \Phi_2^* + \bar{Q}_L^3\, \bar{y}^u T_R \Phi_2 + \bar{Q}_L^3\, \tilde{y}^u u_R \Phi_3 + \bar{Q}_L^{1,2}\, y^d d_R \Phi_3^* \\ + \bar{Q}_L^{1,2}\, \bar{y}_X^d \hat{d}_R \Phi_X + \bar{Q}_L^3\, \bar{y}_X^u T_R \Phi_X + \text{h.c.}\,, \tag{3.30}$$

and

$$\mathcal{L}_Y^\ell = y^\ell \bar{\psi}_L e_R \Phi_1 + y^a \overline{\psi_L^c} \psi_L \Phi_1 + y^s \bar{\psi}_L\, s\, \Phi_2 + \frac{m_s}{2}\, \overline{s^c}\, s + \text{h.c.}\,. \tag{3.31}$$

We defined $\hat{d}_R \equiv (D_R, S_R)$. We note that Eq. (3.31) leads to an inverse seesaw mechanism for neutrino masses [77, 78]. Here, y^a is anti-symmetric while m_s is symmetric, whereas the rest of Yukawa couplings are generic matrices, including those in Eq. (3.30). An additional term $y_X^s \bar{\psi}_L s \Phi_X$ could be added to Eq. (3.31), but given that $\langle \Phi_X \rangle = 0$, it does not contribute to neutrino masses and we will drop it for simplicity. Finally, the scalar potential is given by

$$V = \sum_i \mu_i^2 |\Phi_i|^2 + \lambda_i |\Phi_i|^4 + \sum_{i \neq j} \lambda_{ij} |\Phi_i|^2 |\Phi_j|^2 \\ + f\left(\Phi_1 \Phi_2 \Phi_3 + \text{h.c.}\right) + \frac{\kappa}{2}\left[(\Phi_2^\dagger \Phi_X)^2 + \text{h.c.}\right]\,, \tag{3.32}$$

where $i = 1, 2, 3, X$. The $\mathbb{Z}_2$-soft-breaking term, $f \Phi_1 \Phi_2 \Phi_3$, is required to break unwanted accidental symmetries in the scalar potential.

We will assume the following symmetry breaking pattern

$$\langle \Phi_1 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} k_1 \\ 0 \\ 0 \end{pmatrix}, \quad \langle \Phi_2 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ 0 \\ n \end{pmatrix}, \quad \langle \Phi_3 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ k_3 \\ 0 \end{pmatrix}, \quad \langle \Phi_X \rangle = \begin{pmatrix} 0 \\ 0 \\ 0 \end{pmatrix}. \tag{3.33}$$

3.2.2 331 model with exotic charges

- **Reference:** [33] (see also [74, 75] for similar constructions)
- **Model name:** `331/v2`

Now, we will consider a 331 variant with $\beta = -\sqrt{3}$, as discussed in the context of the diphoton excess in [33]. The fermionic and scalar particle content of the model is summarized in Table 27. In addition, the model contains 17 gauge bosons: the usual 8 gluons; 8 W_i bosons associated to $SU(3)_L$ and the B boson associated to $U(1)_\mathcal{X}$.

The fermionic $SU(3)_L$ triplet representations of the model can be decomposed as

$$\psi_L = \begin{pmatrix} \ell^- \\ -\nu \\ E^{--} \end{pmatrix}_L^{e,\mu,\tau}, \; Q_L^1 = \begin{pmatrix} u \\ d \\ D \end{pmatrix}_L, \; Q_L^2 = \begin{pmatrix} c \\ s \\ S \end{pmatrix}_L, \; Q_L^3 = \begin{pmatrix} b \\ -t \\ T \end{pmatrix}_L. \tag{3.34}$$

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_\mathcal{X}$
ρ	1	$\mathbf{1}$	$\mathbf{3}$	1
η	1	$\mathbf{1}$	$\mathbf{3}$	0
χ	1	$\mathbf{1}$	$\mathbf{3}$	-1
ψ_L	3	$\mathbf{1}$	$\bar{\mathbf{3}}$	-1
e_R	3	$\mathbf{1}$	$\mathbf{1}$	-1
E_R	3	$\mathbf{3}$	$\mathbf{3}$	-2
$Q_L^{1,2}$	2	$\mathbf{3}$	$\mathbf{3}$	$\frac{2}{3}$
Q_L^3	1	$\mathbf{3}$	$\bar{\mathbf{3}}$	$-\frac{1}{3}$
u_R	3	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$
T_R	1	$\mathbf{3}$	$\mathbf{1}$	$-\frac{4}{3}$
d_R	3	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$
D_R, S_R	2	$\mathbf{3}$	$\mathbf{1}$	$\frac{5}{3}$

Table 27. Fermionic and scalar particle content of the 331-v2 model. The scalar and fermion fields are shown in the top and bottom of the table respectively.

Due to the choice $\beta = -\sqrt{3}$, the electric charges for the extra quarks that constitute the third components of the $SU(3)_L$ triplets $Q_L^{1,2,3}$ are 5/3, 5/3 and $-4/3$, respectively. The scalar triplets can be written as

$$\rho = \begin{pmatrix} \rho^+ \\ \rho^0 \\ \rho^{++} \end{pmatrix}, \quad \eta = \begin{pmatrix} \eta^0 \\ \eta_1^- \\ \eta_2^+ \end{pmatrix}, \quad \chi = \begin{pmatrix} \chi^- \\ \chi^{--} \\ \chi^0 \end{pmatrix}. \tag{3.35}$$

Therefore, the particle spectrum of the model contains the exotic quarks in Eq. (3.34), as well as the doubly-charged fermion E^{--} and the scalars ρ^{++} and χ^{--}.

The Yukawa Lagrangian of the model can be split as

$$\mathcal{L}_Y = \mathcal{L}_Y^q + \mathcal{L}_Y^\ell \,, \tag{3.36}$$

where

$$\begin{aligned} \mathcal{L}_Y^q &= y^d\, \overline{Q_L^{1,2}}\, \rho\, d_R + \tilde{y}^d\, \overline{Q_L^3}\, \eta^*\, d_R \\ &+ y^u\, \overline{Q_L^{1,2}}\, \eta\, u_R + \tilde{y}^u\, \overline{Q_L^3}\, \rho^*\, u_R \\ &+ y^J\, \overline{Q_L^{1,2}}\, \chi\, \hat{d}_R + \tilde{y}^J\, \overline{Q_L^3}\, \chi^*\, T_R + \text{h.c.}\,, \end{aligned} \tag{3.37}$$

where we have defined $\hat{d}_R \equiv (D_R, S_R)$, and

$$\mathcal{L}_Y^\ell = y^\ell\, \overline{\psi_L}\, \eta^*\, e_R + y^E\, \overline{\psi_L}\, \chi^*\, E_R + \text{h.c.}\,. \tag{3.38}$$

We note that the exotic fermions E, D, S and T only couple to the χ scalar triplet, and thus only via its vacuum expectation value (VEV) they will acquire masses. Finally, the scalar potential is given by

$$\begin{aligned} V =& \mu_1^2\,|\rho|^2 + \lambda_1|\rho|^4 + \mu_2^2\,|\eta|^2 + \lambda_2|\eta|^4 + \mu_3^2\,|\chi|^2 + \lambda_3|\chi|^4 + \lambda_{12}|\rho|^2|\eta|^2 + \lambda_{13}|\eta|^2|\chi|^2 \\ &+ \lambda_{23}|\eta|^2|\chi|^2 + +\tilde{\lambda}_{12}(\rho^\dagger\eta)(\eta^\dagger\rho) + \tilde{\lambda}_{13}(\rho^\dagger\chi)(\chi^\dagger\rho) + \tilde{\lambda}_{23}(\eta^\dagger\chi)(\chi^\dagger\eta) \\ &+ \sqrt{2}\,f\left(\epsilon_{ijk}\,\rho^i\eta^j\chi^k + \text{h.c.}\right)\,. \end{aligned} \tag{3.39}$$

We will assume the following symmetry breaking pattern

$$\langle\rho\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v_1 \\ 0 \end{pmatrix}\,, \quad \langle\eta\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} v_2 \\ 0 \\ 0 \end{pmatrix}\,, \quad \langle\chi\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ 0 \\ v_3 \end{pmatrix}\,. \tag{3.40}$$

In this case, the non-zero VEV of χ is responsible for the breaking $SU(3)_L \times U(1)_X \to SU(2)_L \times U(1)_Y$. The requirement that this occurs at a scale much above the EW scale then imposes a hierarchy amongst the VEVs, namely that $v_3 \gg v_1, v_2$. Consequently, one of the CP-even scalar states is predominantly from the χ triplet and decouples from the SM. This scalar is then identified as the candidate for the 750 GeV resonance in this model. The decays of this state into two photons proceed via loops involving the heavy fermions, as well as those involving the charged scalars and additional charged vector bosons.

3.3 Other models

Gauged THD model

- **Reference:** [34]
- **Model name:** `GTHDM`

The GTHDM model [79] comes with an additional gauged $SU(2)_H$ symmetry and a $U(1)_X$ symmetry, which is either global or gauged as well. Since the minimal, gauged version suffers from the fact that two massless vector bosons are present, $U(1)_X$ is treated as global symmetry. The scalar and fermion fields are listed in Table 28.

The Lagrangian of the GTHDM contains the SM Lagrangian, extended by the new terms

$$\begin{aligned} \mathcal{L} =& (Y_d\,q\,d\,H^* - Y_d'\,\chi_d\,d\,\Phi + Y_e\,l\,e\,H^* - Y_e'\,\chi_e\,e\,\Phi \\ &+Y_u\,q\,u\,H - Y_u'\,\chi_u\,u\,\Phi^* + Y_\nu\,l\,\nu\,H - Y_\nu'\,\chi_\nu\,\nu\,\Phi^* + \text{h.c.}) \\ &+ \mu_\Delta^2|\Delta_H|^2 - \mu_H^2|H|^2 - \mu_\Phi^2|\Phi|^2 - M_H H^*\Delta_H H + M_\Phi\Phi^*\Delta_H\Phi - \lambda_D\text{Tr}(\Delta_H^\dagger\Delta_H)^2 \\ &- \lambda_{H\Delta}|H|^2\text{Tr}(\Delta_H^\dagger\Delta_H) - \lambda_H|H|^4 - \lambda_{H\Phi}|H|^2|\Phi|^2 - \lambda_{\Phi\Delta}|\Phi|^2\text{Tr}(\Delta_H^\dagger\Delta_H) - \lambda_\Phi|\Phi|^4\,. \end{aligned} \tag{3.41}$$

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$SU(2)_H$	$U(1)_X$
$H = \begin{pmatrix} H_2^c & H_1^+ \\ H_2^0 & H_1^0 \end{pmatrix}$	1	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{2}$	$\mathbf{2}$	1
$\Delta_H = \begin{pmatrix} \delta^0/\sqrt{2} & (\delta^-)^* \\ \delta^- & -\delta^0/\sqrt{2} \end{pmatrix}$	1	$\mathbf{1}$	$\mathbf{1}$	0	$\mathbf{3}$	0
$\Phi = (\phi^c\,\phi^0)^T$	1	$\mathbf{1}$	$\mathbf{1}$	0	$\mathbf{2}$	-1
$q = (u_L\,d_L)^T$	3	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$	$\mathbf{1}$	0
$l = (\nu_L\,e_L)^T$	3	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	$\mathbf{1}$	0
$d = (d_R^H\,d_R)^T$	3	$\overline{\mathbf{3}}$	$\mathbf{1}$	$\frac{1}{3}$	$\mathbf{2}$	1
$u = (u_R\,u_R^H)^T$	3	$\overline{\mathbf{3}}$	$\mathbf{1}$	$-\frac{2}{3}$	$\mathbf{2}$	-1
$\nu = (\nu_R\,\nu_R^H)^T$	3	$\mathbf{1}$	$\mathbf{1}$	0	$\mathbf{2}$	-1
$e = (e_R^H\,e_R)^R$	3	$\mathbf{1}$	$\mathbf{1}$	1	$\mathbf{2}$	1
χ_d	3	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	$\mathbf{1}$	0
χ_u	3	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	$\mathbf{1}$	0
χ_ν	3	$\mathbf{1}$	$\mathbf{1}$	0	$\mathbf{1}$	0
χ_e	3	$\mathbf{1}$	$\mathbf{1}$	-1	$\mathbf{1}$	0

Table 28. Scalar and fermion fields in the `GTHDM`

The breaking of $SU(2)_L \times U(1)_Y \times SU(2)_H \to U(1)_{em}$ is triggered by

$$\langle\Phi\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v_\phi \end{pmatrix}\,, \quad \langle H\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 & 0 \\ 0 & v \end{pmatrix}\,, \quad \langle\Delta_H\rangle = \frac{1}{2}\begin{pmatrix} v_T & 0 \\ 0 & -v_T \end{pmatrix}\,. \tag{3.42}$$

After EWSB, there are three neutral gauge bosons which mix giving rise to the γ, Z, Z' mass eigenstates and two charged ones (W, W') which do not mix. The neutral component of the SM-singlet Φ, ϕ^0, is considered to be the candidate for the 750 GeV resonance while its VEV gives mass to the exotic fermions that are needed to run in the loops. As ϕ^0 will typically mix with the $SU(2)_L$ doublet Higgs, this mixing needs to be suppressed by a specific parameter choice in order to avoid the tight bounds from the dijet, ZZ or dilepton channels.

4 Supersymmetric models

There are several ideas to explain the diphoton excess within a supersymmetric framework. Some of them make use of SUSY models which already exist in the literature, and for which also `SARAH` model files exist: the MSSM with trilinear R-parity violation [80, 81], the simplest models with Dirac gauginos [82], or the model with gauged

$U(1)_L \times U(1)_B$ [83]. We will not make any further comment on these models, but concentrate in the following on the models which are newly implemented.

4.1 NMSSM extensions with vector-like multiplets

- **Reference:** [35–38]
- **Model name:** `NMSSM+VL`

The scalar component of the gauge-singlet superfield $\hat{S}$ of the Next to Minimal Supersymmetric Standard Model (NMSSM) can explain the 750 GeV resonance, if one adds vector-like $SU(5)$ multiplets to enhance the diphoton rate. The new multiplets are added in pairs of $(\mathbf{5}, \bar{\mathbf{5}})$ and/or $(\mathbf{10}, \bar{\mathbf{10}})$ in order to preserve gauge coupling unification. Typically one also imposes a $\mathbb{Z}_2$ symmetry to forbid mixing of the new vector-like particles with the MSSM particles. The authors of Ref. [35] mention the possibility to interpret the resonance as two nearly degenerate singlet-like bosons, roughly the scalar and pseudoscalar components of the singlet $\hat{S}$. There are some differences in the singlet and Higgs superpotential interactions included in the different papers:

- in Ref. [37] the authors assume $\lambda \hat{S} \hat{H}_u \hat{H}_d$ and $\kappa/3 \hat{S}^3$ to be present, but $\mu \hat{H}_u \hat{H}_d$ to be absent;
- in Ref. [36] $\mu \hat{H}_u H_d$, $\lambda \hat{S} \hat{H}_u \hat{H}_d$ and $M_S/2\, \hat{S}^2$ are present;
- in Ref. [38] only $M_S/2\, \hat{S}^2$ is present. This does not cause a mixing between the singlet and Higgs doublet at the tree level, but such a mixing is unavoidable at the loop level.

The `SARAH` implementations use the most general version of the superpotential: all possible interactions are present. The different limits according to the proposals of Refs. [36–38] can be obtained by setting the corresponding parameters to zero in numerical studies. In what follows we describe the models with the vector-like multiplets in different representations of $SU(5)$.

4.2 NMSSM with vectorlike top

- **Model name:** `NMSSM+VL/VLtop`
- **Reference:** [35]

This model is an extension of the NMSSM by a vector-like top. There is a global $\mathbb{Z}_2$ R-parity and a $\mathbb{Z}_3$ symmetry, under which all particles transform as $X \to \exp(i\frac{2\pi}{3})\,X$. In this way, only terms with three superfields are allowed in the superpotential. The particle content is given in Table 29. Since this model does not introduce complete multiplets (**5** or **10**) of $SU(5)$, gauge coupling unification is not achieved. The superpotential is given by

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$\mathbb{Z}_2$	$\mathbb{Z}_3$
$\hat{S}$	1	$\mathbf{1}$	$\mathbf{1}$	0	$+$	$\exp(i\frac{2\pi}{3})$
$\hat{T}$	1	$\bar{\mathbf{3}}$	$\mathbf{1}$	$-\frac{2}{3}$	$-$	$\exp(i\frac{2\pi}{3})$
$\hat{T}'$	1	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	$-$	$\exp(i\frac{2\pi}{3})$

Table 29. Superfield content beyond the MSSM superfields, including a singlet and a vector-like top.

$$W = -\,Y_d\,\hat{d}\,\hat{q}\hat{H}_d\,-Y_e\,\hat{e}\,\hat{l}\hat{H}_d\,+Y_u\,\hat{u}\,\hat{q}\hat{H}_u + \frac{1}{3}\kappa\hat{S}^3+\lambda\hat{S}\hat{H}_u\hat{H}_d+Y_t\hat{T}\hat{Q}\hat{H}_u+\lambda_T\hat{T}\hat{T}'\hat{S}+\lambda_U\hat{U}\hat{T}'\hat{S} \tag{4.1}$$

Beyond the neutral scalar components of the two Higgs doublets, after EWSB the complex singlet gets a VEV and can be decomposed as

$$S = \frac{1}{\sqrt{2}}\left(v_S+\phi_S+i\,\sigma_S\right). \tag{4.2}$$

The fermionic components of $\hat{T}, \hat{T}'$ mix with the up-type quarks, while the scalar components mix with the up-like squarks.

4.2.1 Pairs of 5 of $SU(5)$

- **Model name:** `NMSSM+VL/5plets`

The superfields beyond the MSSM are shown in Table 30. In the current implementation we only have one copy of $(\mathbf{5},\bar{\mathbf{5}})$ fields, but having at least three copies of them should give a better fit to the diphoton resonance. According to [36] the fit is even better with four copies, however in that case one might hit a Landau pole.

The superpotential is given by

$$\begin{aligned} W =&-\,Y_d\,\hat{d}\,\hat{q}\hat{H}_d\,-Y_e\,\hat{e}\,\hat{l}\hat{H}_d\,+Y_u\,\hat{u}\,\hat{q}\hat{H}_u+\mu\hat{H}_u\hat{H}_d\\ &+\frac{1}{3}\kappa\hat{S}^3+\lambda\hat{S}\hat{H}_u\hat{H}_d+M_S\hat{S}^2+t_S\hat{S}\\ &+\lambda_D\hat{S}\hat{D}\hat{D}'+\lambda_L\hat{S}\hat{L}\hat{L}'+M_L\hat{L}\hat{L}'+M_D\hat{D}\hat{D}'\,. \end{aligned} \tag{4.3}$$

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$
$\hat{S}$	1	$\mathbf{1}$	$\mathbf{1}$	0
$\hat{D}$	1	$\overline{\mathbf{3}}$	$\mathbf{1}$	$\frac{1}{3}$
$\hat{D}'$	1	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$
$\hat{L}$	1	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$
$\hat{L}'$	1	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{2}$

Table 30. Superfield content in the case of a pair of **5**'s of $SU(5)$.

Beyond the neutral scalar components of the two Higgs doublets, also the singlet gets a VEV after EWSB and can be decomposed as in Eq. (4.2).

4.3 Pairs of 10 of $SU(5)$

- **Model name:** `NMSSM+VL/10plets`

The superfields beyond the MSSM are shown in Table 31. The superpotential is

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$
$\hat{S}$	1	$\mathbf{1}$	$\mathbf{1}$	0
$\hat{U}$	1	$\overline{\mathbf{3}}$	$\mathbf{1}$	$-\frac{2}{3}$
$\hat{U}'$	1	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$
$\hat{Q}$	1	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$
$\hat{Q}'$	1	$\overline{\mathbf{3}}$	$\mathbf{2}$	$-\frac{1}{6}$
$\hat{E}$	1	$\mathbf{1}$	$\mathbf{1}$	1
$\hat{E}'$	1	$\mathbf{1}$	$\mathbf{1}$	-1

Table 31. Superfield content in the case of a pair of **10**'s of $SU(5)$.

given by

$$\begin{aligned} W = &- Y_d\, \hat{d}\, \hat{q} \hat{H}_d \; - Y_e\, \hat{e}\, \hat{l} \hat{H}_d \; + Y_u\, \hat{u}\, \hat{q} \hat{H}_u + \mu \hat{H}_u \hat{H}_d \\ &+ \frac{1}{3}\kappa \hat{S}^3 + \lambda \hat{S}\hat{H}_u\hat{H}_d + M_S \hat{S}^2 + t_S \hat{S} \\ &+ Y_{10}\hat{Q}\hat{U}\hat{H}_u + Y'_{10}\hat{Q}'\hat{U}'\hat{H}_d + \lambda_Q \hat{S}\hat{Q}\hat{Q}' + \lambda_U \hat{S}\hat{U}\hat{U}' + \lambda_E \hat{S}\hat{E}\hat{E}' \\ &+ M_U \hat{U}\hat{U}' + M_Q \hat{Q}\hat{Q}' + M_E \hat{E}\hat{E}' \,. \end{aligned} \tag{4.4}$$

The symmetry breaking pattern is the same as for the model with 5-plets.

4.4 Pairs of 5 and 10 of $SU(5)$

- **Model name:** `NMSSM+VL/5+10plets`

This model combines the previous two setups, adding pairs of vectorlike **5** and **10** representations of $SU(5)$. The superfields beyond the MSSM are shown in Table 32. The superpotential is given by

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$
$\hat{S}$	1	$\mathbf{1}$	$\mathbf{1}$	0
$\hat{D}$	1	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\frac{1}{3}$
$\hat{D}'$	1	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$
$\hat{L}$	1	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$
$\hat{L}'$	1	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{2}$
$\hat{U}$	1	$\bar{\mathbf{3}}$	$\mathbf{1}$	$-\frac{2}{3}$
$\hat{U}'$	1	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$
$\hat{Q}$	1	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$
$\hat{Q}'$	1	$\bar{\mathbf{3}}$	$\mathbf{2}$	$-\frac{1}{6}$
$\hat{E}$	1	$\mathbf{1}$	$\mathbf{1}$	1
$\hat{E}'$	1	$\mathbf{1}$	$\mathbf{1}$	-1

Table 32. Superfield content in the case of a pair of **5**'s and **10**'s of $SU(5)$.

$$\begin{aligned} W = &- Y_d\, \hat{d}\, \hat{q} \hat{H}_d \, - Y_e\, \hat{e}\, \hat{l} \hat{H}_d \, + Y_u\, \hat{u}\, \hat{q} \hat{H}_u + \mu \hat{H}_u \hat{H}_d \\ &+ \frac{1}{3}\kappa \hat{S}^3 + \lambda \hat{S} \hat{H}_u \hat{H}_d + M_S \hat{S}^2 + t_S \hat{S} \\ &+ Y'_D\, \hat{D}\, \hat{Q} \hat{H}_d \, + Y'_E\, \hat{E}\, \hat{L} \hat{H}_d \, + Y'_U\, \hat{U}\, \hat{Q} \hat{H}_u + Y''_D\, \hat{D}'\, \hat{Q}' \hat{H}_u \, + Y''_E\, \hat{E}'\, \hat{L}' \hat{H}_u \, + Y''_U\, \hat{U}'\, \hat{Q}' \hat{H}_d \\ &+ \lambda_D \hat{S} \hat{D} \hat{D}' + \lambda_L \hat{S} \hat{L} \hat{L}' + \lambda_Q \hat{S} \hat{Q} \hat{Q}' + \lambda_E \hat{S} \hat{E} \hat{E}' + \lambda_U \hat{S} \hat{U} \hat{U}' \\ &+ M_L \hat{L} \hat{L}' + M_D \hat{D} \hat{D}' + M_Q \hat{Q} \hat{Q}' + M_E \hat{E} \hat{E}' + M_U \hat{U} \hat{U}'\,. \end{aligned} \tag{4.5}$$

The symmetry breaking pattern is the same as for the model with 5-plets.

4.5 Pairs of 5 of $SU(5)$ and R-parity violation

- **Model name:** `NMSSM+VL/5plets+RpV`

One can relax the assumption of the $\mathbb{Z}_2$ symmetry that forbids mixing between vector-like fields and MSSM fields, in which case terms like $\kappa_5 \hat{S} \hat{L} \hat{H}_u$ are added to model [36]. Furthermore, this also breaks R-parity.

The superfields beyond the MSSM are shown in Table 30 and the superpotential is given by

$$\begin{aligned} W = & -Y_d\,\hat{d}\,\hat{q}\hat{H}_d\; - Y_e\,\hat{e}\,\hat{l}\hat{H}_d\; + Y_u\,\hat{u}\,\hat{q}\hat{H}_u + \mu\hat{H}_u\hat{H}_d \\ & + \frac{1}{3}\kappa\hat{S}^3 + \lambda\hat{S}\hat{H}_u\hat{H}_d + M_S\hat{S}^2 + t_S\hat{S} \\ & + \kappa_5\hat{S}\hat{L}\hat{H}_u + \kappa_5'\hat{S}\hat{L}'\hat{H}_d + \lambda_D\hat{S}\hat{D}\hat{D}' + \lambda_L\hat{S}\hat{L}\hat{L}' + M_L\hat{L}\hat{L}' + M_D\hat{D}\hat{D}'\,. \end{aligned} \tag{4.6}$$

The inclusion of R-parity violating terms triggers VEVs for the neutral components of $\tilde{L}$ and $\tilde{L}'$,

$$\langle\tilde{L}\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v_L \end{pmatrix}\,, \quad \langle\tilde{L}'\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix} 0 \\ v_{L'} \end{pmatrix}\,, \tag{4.7}$$

and causes mixing between the vector-like states and the Higgs components.

4.6 Broken MRSSM

- **Reference:** [39]
- **Model name:** `brokenMRSSM`

In the minimal R-supersymmetric model (MRSSM) the scalar R_u (see Table 33) is proposed as an explanation to the 750 GeV resonance. In order to explain the diphoton excess, it is necessary to add explicitly an R-symmetry breaking term to the Lagrangian

$$L_{\not{R}} = T_u H_u \tilde{Q}\tilde{u}\,, \tag{4.8}$$

where T_u is a dimensionful trilinear coupling. This source of R-symmetry breaking has several consequences, not discussed in Ref. [39], which however are taken into account in the model implementation:

1. The term in Eq. (4.8) will introduce Majorana gaugino masses via RGE effects
2. The Majorana gaugino masses will also generate all other trilinear and bilinear soft-terms
3. This causes R-symmetry breaking terms R_iH_i $i = d, u$ which will trigger VEVs for the R-fields
4. The neutralinos and gluinos are no longer Dirac particles, but mix to Majorana fermions
5. There is a mixing between fields of different R-charges.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$
$\hat{S}$	1	$\mathbf{1}$	$\mathbf{1}$	0
$\hat{T}$	1	$\mathbf{1}$	$\mathbf{3}$	0
$\hat{O}$	1	$\mathbf{8}$	$\mathbf{1}$	0
$\hat{R}_d$	1	$\mathbf{1}$	$\mathbf{2}$	$+\frac{1}{2}$
$\hat{R}_u$	1	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$

Table 33. Superfields of the broken MRSSM beyond the MSSM particle content.

The superfields beyond the MSSM are listed in Table 33. The superpotential, assumed to conserve R-symmetry, is given by

$$
\begin{aligned}
W = &- Y_d\,\hat{d}\,\hat{q}\hat{H}_d \,-Y_e\,\hat{e}\,\hat{l}\hat{H}_d \,+Y_u\,\hat{u}\,\hat{q}\hat{H}_u + \mu_D\,\hat{R}_d\hat{H}_d \\
&+ \mu_U\,\hat{R}_u\hat{H}_u \,+\hat{S}(\lambda_d\,\hat{R}_d\hat{H}_d \,+\lambda_u\,\hat{R}_u\hat{H}_u) \\
&+ \lambda_d^T\,\hat{R}_d\hat{T}\,\hat{H}_d \,+\lambda_u^T\,\hat{R}_u\hat{T}\,\hat{H}_u \ .
\end{aligned}
\tag{4.9}
$$

As explained above, because the R-symmetry is broken in the soft sector of the model, all possible tri- and bilinear soft-breaking terms corresponding to the superpotential terms will be generated.

The following VEVs appear after EWSB, beyond those of the neutral scalar components of the two Higgs doublets

$$
\langle R_d\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix}0\\ v_{R_d}\end{pmatrix}, \quad \langle R_u\rangle = \frac{1}{\sqrt{2}}\begin{pmatrix}0\\ v_{R_u}\end{pmatrix}, \quad \langle S\rangle = \frac{v_S}{\sqrt{2}}, \quad \langle T\rangle = \frac{1}{2}\begin{pmatrix}v_T & 0\\ 0 & -v_T\end{pmatrix} .
\tag{4.10}
$$

The authors favor to have large stop mixing for a not too large R-symmetry breaking term T_u by considering the limit $v_d \sim v_u$ and $m_{\tilde{t}_L} \sim m_{\tilde{t}_R}$. However, in this limit the mass of the SM-like Higgs is tiny and often tachyonic: in the MSSM, the Higgs tree-level mass vanishes for $\tan\beta \to 1$, and this model has additional negative contributions to the mass from the new D-terms present in models with Dirac gauginos. It is also questionable if the case with very large T_u is a viable scenario because these values are highly restricted by charge and colour breaking minima [84, 85], which demands careful checks. This is similar to the vacuum stability issues discussed in [1].

4.7 $U(1)'$-extended MSSM

- **Reference:** [40, 41]
- **Model name:** `MSSM+U1prime-VL`

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)'$
$\hat{Q}$	3	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$	$\frac{1}{2}$
$\hat{d}^c$	3	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\frac{1}{3}$	$\frac{1}{2}$
$\hat{u}^c$	3	$\bar{\mathbf{3}}$	$\mathbf{1}$	$-\frac{2}{3}$	$\frac{1}{2}$
$\hat{L}$	3	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	$\frac{1}{2}$
$\hat{e}^c$	3	$\mathbf{1}$	$\mathbf{1}$	1	$\frac{1}{2}$
$\hat{\nu}^c$	3	$\mathbf{1}$	$\mathbf{1}$	0	$\frac{1}{2}$
$\hat{H}_d$	1	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	-1
$\hat{H}_u$	1	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{2}$	-1

Table 34. Quantum numbers of the MSSM fields under the full gauge group in the `MSSM+U1prime-VL`.

Field	Gen.	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)'$
$\hat{T}$	2	$\mathbf{3}$	$\mathbf{1}$	$\frac{2}{3}$	-1
$\hat{T}^c$	2	$\bar{\mathbf{3}}$	$\mathbf{1}$	$-\frac{2}{3}$	-1
$\hat{T}_3$	1	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	-1
$\hat{T}_3^c$	1	$\bar{\mathbf{3}}$	$\mathbf{1}$	$\frac{1}{3}$	-1
$\hat{D}$	2	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{2}$	-1
$\hat{D}^c$	2	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	-1
$\hat{X}$	1	$\mathbf{1}$	$\mathbf{1}$	1	2
$\hat{X}^c$	1	$\mathbf{1}$	$\mathbf{1}$	-1	2
$\hat{N}$	1	$\mathbf{1}$	$\mathbf{1}$	0	2
$\hat{N}^c$	1	$\mathbf{1}$	$\mathbf{1}$	0	2
S	1	$\mathbf{1}$	$\mathbf{1}$	0	2
S^c	1	$\mathbf{1}$	$\mathbf{1}$	0	-2
S_1	1	$\mathbf{1}$	$\mathbf{1}$	0	-4
S_1^c	1	$\mathbf{1}$	$\mathbf{1}$	0	4
S_2	1	$\mathbf{1}$	$\mathbf{1}$	0	-2

Table 35. Extra superfield content of the `MSSM+U1prime-VL` and their quantum numbers under the full gauge group.

In this model all MSSM fields carry a non-zero $U(1)'$-charge so that anomaly cancellation requires additional superfields (see Table 34), which are also responsible for the spontaneous $U(1)'$ breaking. Furthermore, colour-charged and colour-uncharged matter superfields which are vector-like with respect to the MSSM gauge group are introduced. A combination of scalar singlets S and S_i is supposed to give the 750 GeV resonance.

The complete superfield content with all gauge quantum numbers is given in Tables 34 and 35. In addition to the usual matter parity, we impose a $\mathbb{Z}_2$ symmetry under which all exotic matter superfields are odd and all other superfields are even in order to reduce the number of superpotential terms and hence reduce the complexity of the model.

The superpotential is given by

$$\begin{aligned} W = & - Y_d\, \hat{d}^c\, \hat{Q}\, \hat{H}_d\, - Y_e\, \hat{e}^c\, \hat{L}\, \hat{H}_d\, + Y_u\, \hat{u}^c\, \hat{Q}\, \hat{H}_u + Y_\nu\, \hat{\nu}^c\, \hat{L}\, \hat{H}_u + \lambda\, \hat{S}\, \hat{H}_u\, \hat{H}_d \\ & + \lambda_N\, \hat{S}_1\, \hat{N}\, \hat{N}^c + \lambda_D\, \hat{S}\, \hat{D}\, \hat{D}^c + \lambda_X\, \hat{S}_1\, \hat{X}\, \hat{X}_c + \lambda_T\, \hat{S}\, \hat{T}^c\, \hat{T} + \lambda_{T3}\, \hat{S}\, \hat{T}_3^c\, \hat{T}_3 \\ & + \mu_S\, \hat{S}\, \hat{S}^c + \mu_{1S}\, \hat{S}_1\, \hat{S}_1^c + \mu_{2S}\, \hat{S}\, \hat{S}_2 + \kappa_1\, \hat{S}\, \hat{S}\, \hat{S}_1 + \kappa_2\, \hat{S}^c\, \hat{S}_2\, \hat{S}_1^c\,. \end{aligned} \tag{4.11}$$

In addition to the neutral components of the two Higgs doublets, the MSSM singlets get VEVs according to

$$\langle S_i^{(c)} \rangle = \frac{v_{S_i}^{(c)}}{\sqrt{2}}\,. \tag{4.12}$$

4.8 E_6-inspired SUSY models with extra $U(1)$ model

- **Reference:** [42]
- **Model name:** `SUSYmodels/E6SSMalt`

E_6-inspired SUSY models predict extra SM-gauge singlets and extra exotic fermions, so they immediately have the ingredients that many authors have tried to use to fit the diphoton excess. These models are often motivated as a solution to the μ-problem of the MSSM, because the extra $U(1)$ gauge symmetry forbids the μ-term, while when one of the singlet fields develops a VEV at the TeV scale this breaks the extra $U(1)$ giving rise to a massive Z' vector boson and at the same time generates an effective μ term through the singlet interaction with the up- and down-type Higgs fields, $\lambda \hat{S}\hat{H}_u\hat{H}_d$. The matter content of the model at low energies fills three generations of complete **27**-plet representations of E_6, which ensures that anomalies automatically cancel.

A number of models of this nature have been proposed as explanations of the diphoton excess [42, 86, 87]. The example we implement here [42] is a variant of the E_6SSM [88, 89]. In this version two singlet states develop VEVs and the idea is that the 750 GeV excess is explained by one of these singlet states with a loop-induced decay through the exotic states.

In E_6 models the extra $U(1)$ which extends the SM gauge group is given as a linear combination of $U(1)_\psi$ and $U(1)_\chi$ which appear from the breakdown of the E_6 symmetry as $E_6 \to SO(10) \times U(1)_\psi$ followed by $SO(10)$ into $SU(5)$, $SO(10) \to SU(5) \times U(1)_\chi$. In the E_6SSM and the variant implemented here the specific combination is,

$$U(1)_N = \frac{1}{4}U(1)_\chi + \frac{\sqrt{15}}{4}U(1)_\psi. \tag{4.13}$$

To allow one-step gauge coupling unification however some incomplete multiplets must be included in the low energy matter content. So in addition to the matter filling complete **27** representations of E_6 there are also two $SU(2)$ multiplets $\hat{H}'$ and $\hat{\overline{H}}'$, which are the only components from additional $\mathbf{27'}$ and $\overline{\mathbf{27'}}$ that survive to low energies. All gauge anomalies cancel between these two states, so they do not introduce any gauge anomalies. Furthermore, the low energy matter content of the model beyond the MSSM one includes three generations of exotic diquarks[7], $\hat{D}_i$, $\hat{\overline{D}}_i$, three generations of SM singlet superfields $\hat{S}_i$ and extra Higgs-like states $H^u_{1,2}$ and $H^d_{1,2}$ that do not get VEVs.

The full set of superfields are given in Table 36 along with their representations under $SU(3)$ and $SU(2)$ and the charges of the two U(1) gauge groups and the discrete symmetries, which we will now discuss.

The $\mathbb{Z}_2^L$ symmetry plays a role similar to R-parity in the MSSM, being imposed to avoid rapid proton decay in the model. However with this imposed there are still terms in the superpotential that can lead to dangerous flavour changing neutral currents (FCNCs). To forbid these, an approximate $\mathbb{Z}_2^H$ symmetry is imposed. In the original E_6SSM model only $\hat{S}_3$, $\hat{H}_d$ and $\hat{H}_u$ were even under the $\mathbb{Z}_2^H$ symmetry, however in this variant S_2 is also even under this approximate symmetry.

The full superpotential before imposing any discrete symmetries is given by

$$W_{E6} = W_0 + W_1 + W_2, \tag{4.14}$$

where

$$\begin{aligned} W_0 &= \lambda_{ijk}\hat{S}_i\hat{H}^d_j\hat{H}^u_k + \kappa_{ijk}\hat{S}_i\hat{D}_j\hat{\overline{D}}_k + h^N_{ijk}\hat{N}^c_i\hat{H}^u_j\hat{L}_k \\ &\quad + h^U_{ijk}\hat{u}^c_i\hat{H}^u_j\hat{Q}_k + h^D_{ijk}\hat{d}^c_i\hat{H}^d_j\hat{Q}_k + h^E_{ijk}\hat{e}^c_i\hat{H}^d_j\hat{L}_k, \end{aligned} \tag{4.15}$$

$$W_1 = g^Q_{ijk}\hat{D}_i\hat{Q}_j\hat{Q}_k + g^q_{ijk}\hat{\overline{D}}_i\hat{d}^c_j\hat{u}^c_k, \tag{4.16}$$

$$W_2 = g^N_{ijk}\hat{N}^c_i\hat{D}_j\hat{d}^c_k + g^E_{ijk}\hat{e}^c_i\hat{D}_j\hat{u}^c_k + g^D_{ijk}\hat{Q}_i\hat{L}_j\hat{\overline{D}}_k. \tag{4.17}$$

[7] In the original E_6SSM these states could be either diquark or leptoquark in nature, depending on the choice of a discrete symmetry, but in the model considered here the allowed superpotential terms for the decay of these exotic quarks imply they are diquark.

Field	Gen	$SU(3)_C$	$SU(2)_L$	$U(1)_Y$	$U(1)_N$	$\mathbb{Z}_2^H$	$\mathbb{Z}_2^L$
$\hat{Q}_i$	3	$\mathbf{3}$	$\mathbf{2}$	$\frac{1}{6}$	1	-	+
$\hat{u}_i^c$	3	$\overline{\mathbf{3}}$	$\mathbf{1}$	$-\frac{2}{3}$	1	-	+
$\hat{d}_i^c$	3	$\overline{\mathbf{3}}$	$\mathbf{1}$	$\frac{1}{3}$	2	-	+
$\hat{L}_i$	3	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	2	-	-
$\hat{e}_i^c$	3	$\mathbf{1}$	$\mathbf{1}$	1	1	-	-
$\hat{N}_i^c$	3	$\mathbf{1}$	$\mathbf{1}$	0	0	-	-
$\hat{S}_i$	2	$\mathbf{1}$	$\mathbf{1}$	0	5	+	+
$\hat{S}_1$	1	$\mathbf{1}$	$\mathbf{1}$	0	5	-	+
$\hat{H}_u$	1	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{2}$	-2	+	+
$\hat{H}_d$	1	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	-3	+	+
$\hat{H}_\alpha^u$	2	$\mathbf{1}$	$\mathbf{2}$	$\frac{1}{2}$	-2	-	+
$\hat{H}_\alpha^d$	2	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	-3	-	+
$\hat{D}_i$	3	$\mathbf{3}$	$\mathbf{1}$	$-\frac{1}{3}$	-2	-	+
$\hat{\overline{D}}$	3	$\overline{\mathbf{3}}$	$\mathbf{1}$	$\frac{1}{3}$	-3	-	+
$\hat{L}_4$	1	$\mathbf{1}$	$\mathbf{2}$	$-\frac{1}{2}$	2	-	+
$\hat{\overline{L}}_4$	1	$\mathbf{1}$	$\overline{\mathbf{2}}$	$\frac{1}{2}$	-2	-	+

Table 36. The representations of the chiral superfields under the $SU(3)_C$ and $SU(2)_L$ gauge groups, and their $U(1)_Y$ and $U(1)_N$ charges without the E_6 normalisation. The GUT normalisations are $\sqrt{\frac{5}{3}}$ for $U(1)_Y$ and $\sqrt{40}$ for $U(1)_N$. The transformation properties under the discrete symmetries $\mathbb{Z}_2^H$, $\mathbb{Z}_2^L$ are also shown, where '+' indicates the superfield is even under the symmetry and '−' indicates that it is odd under the symmetry.

However, with the discrete symmetries imposed and integrating out the heavy right-handed neutrinos, the superpotential in this specific variant reduces to[8],

$$W_{E_6SSM\,\text{variant}} = W_{MSSM}^{(\mu=0)} + \sum_{\alpha=2}^{3}\sum_{i=1}^{3}\hat{S}^\alpha(\lambda_{\alpha\, i}\hat{H}_u^i\hat{H}_d^i + \kappa_{\alpha\, i}\hat{D}^i\hat{\overline{D}}^i)$$
$$+\mu'\hat{H}\prime\hat{\overline{H}}' + h_{4j}^E(\hat{H}_d\hat{H}')\hat{e}_j^c \tag{4.18}$$

One should remember that the $\mathbb{Z}_2^H$ can only be an approximate symmetry as otherwise the exotic quarks could not decay. In this variant the exotic quarks decay therough the $\mathbb{Z}_2^H$ violating interactions of W_1.

[8] In the paper proposing this variant to explain the excess [42], the terms involving the surviving Higgs states on the second line are omitted from the superpotential.

In the paper it is assumed that the singlet mixing can be negligible and the numerical calculation was performed under this assumption, neglecting any mixing between the singlet state which decays to $\gamma\gamma$ via the exotic states and the other CP-even Higgs states from the standard $SU(2)$ doublets. However it is clear that there must be some mixing from the D-terms, and therefore if that is included one important check would be to test whether other decays are sufficiently suppressed. Moreover, the parameters needed to simultaneously get a 125 GeV SM-like Higgs state and a 750 GeV singlet-dominated state are not given. In this respect we note that the singlet VEVs appear both in the diagonal entries of the mass matrix and in the off-diagonal entries that mix the singlet states with the doublet states.

We finally note that other similar E_6 models have also been proposed in the context of the diphoton excess. These include a model by two authors from the original paper [87], a model with a different $U(1)$ group at low energies [90], and a model that is still E_6-inspired, but where no extra $U(1)$ survives down to low energies [91].

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