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# Enhancement of $h \to \gamma\gamma$ by seesaw-motivated exotic scalars

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ABSTRACT

We examine the role of seesaw-motivated exotic scalars in loop-mediated Higgs decays. We consider a simple TeV-scale seesaw model built upon the fermionic quintuplet mediator in conjunction with the scalar quadruplet, where we examine portions of the model parameter space for which the contributions of charged components of the scalar quadruplet significantly increase the $h \to \gamma\gamma$ decay rate. The most significant change in the diphoton width comes from a doubly charged scalar $\Phi^{--}$ which should be the lightest component in the scalar quadruplet. In the part of the parameter space where the $h \to \gamma\gamma$ decay width is enhanced by a factor 1.25–2 there is a mild suppression of the $h \to Z\gamma$ decay width by a factor 0.9–0.7.



## 1. Introduction

The recently discovered resonance with mass $m_h \simeq$ 125–126 GeV [1,2] strongly resembles the Standard Model (SM) Higgs boson. There is a hint that the loop-induced $h \to \gamma\gamma$ event rate [3,4] deviates, modulo QCD uncertainties [5], by a factor 1.5–2 [6,7] from its SM value,

$$\frac{[\sigma(gg \to h) \times \mathrm{BR}(h \to \gamma\gamma)]_{\mathrm{LHC}}}{[\sigma(gg \to h) \times \mathrm{BR}(h \to \gamma\gamma)]_{\mathrm{SM}}} = 1.71 \pm 0.33. \tag{1}$$

This indicates the existence of additional charged particle(s) on which the tree-level Higgs decay modes are much less sensitive. The history of the "prediscovery" of charmed and top quarks through the loop amplitudes may be repeated in the scalar sector.

On the other hand, if this enhancement in $h \to \gamma\gamma$ disappears with a larger integrated luminosity, it will still constrain the parameter space of various extensions of the SM containing new charged states which should affect this loop amplitude. In this spirit there is a number of theoretical attempts to match the indicated discrepancy by the effects of an extended Higgs sector [8–11].

Notable, the $h \to \gamma\gamma$ decay rate is sensitive on the doubly charged scalar fields such as those existing in Higgs triplet model of neutrino mass generation. These states have been in focus of both the direct searches at the LHC [12,13] and of theoretical investigations [14–18]. Here we go a step further by studying the effects of singly and doubly charged components of a scalar quadruplet contained in our recent TeV-scale seesaw model for neutrino masses [19].

The Letter is organized as follows. In Section 2 we briefly review the TeV-scale seesaw model at hand. A detailed study of the scalar potential is presented in Section 3. We investigate the effects of exotic scalars on the loop-level Higgs decays in Sections 4 and 5. Possible direct bounds on exotic scalars are discussed in Section 6. The conclusion is presented in Section 7.

## 2. Fermionic quintuplet seesaw model

A simple and predictive TeV-scale seesaw model [19] under consideration belongs to a class of beyond dimension-five seesaw models [20–22] which, on account of introducing higher isomultiplets, lower the originally high seesaw scale to the TeV scale accessible at the LHC. In order to produce the seesaw diagram in the simple model at hand, the fermionic hypercharge-zero quintuplets $\Sigma_R = (\Sigma_R^{++}, \Sigma_R^{+}, \Sigma_R^{0}, \Sigma_R^{-}, \Sigma_R^{--})$ transforming as $(1, 5, 0)$ under the SM gauge group have to be accompanied by a scalar quadruplet $\Phi = (\Phi^{+}, \Phi^{0}, \Phi^{-}, \Phi^{--})$ transforming as $(1, 4, -1)$.

The gauge invariant and renormalizable Lagrangian involving these new fields reads

$$\mathcal{L} = \overline{\Sigma_R} i\gamma^\mu D_\mu \Sigma_R + \left(D^\mu \Phi\right)^\dagger (D_\mu \Phi) - \left(\overline{L_L} Y \Phi \Sigma_R + \frac{1}{2}\overline{(\Sigma_R)^C} M \Sigma_R + \mathrm{H.c.}\right) - V(H, \Phi), \tag{2}$$

where the scalar potential contains the mass and the quartic terms for the doublet $H$ and the quadruplet $\Phi$ fields

$$V(H, \Phi) = -\mu_H^2 H^\dagger H + \mu_\Phi^2 \Phi^\dagger \Phi + \lambda_1 \left(H^\dagger H\right)^2 + \lambda_2 H^\dagger H \Phi^\dagger \Phi$$

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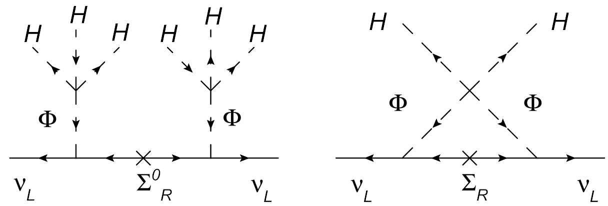


**Fig. 1.** Tree-level dimension-nine operator and one-loop dimension-five operator diagrams relevant for generating the light neutrino masses.

$$+ \lambda_3 H^* H \Phi^* \Phi + (\lambda_4 H^* H H \Phi + \text{H.c.}) + (\lambda_5 H H \Phi \Phi + \text{H.c.}) + (\lambda_6 H \Phi^* \Phi \Phi + \text{H.c.}) + \lambda_7 (\Phi^\dagger \Phi)^2 + \lambda_8 \Phi^* \Phi \Phi^* \Phi. \tag{3}$$

In the tensor notation the terms in Eqs. (2) and (3) read [19]

$$\begin{aligned}
\overline{L_L} \Phi \Sigma_R &= \overline{L_L}^i \Phi_{jkl} \Sigma_{Rij'k'l'} \epsilon^{jj'} \epsilon^{kk'} \epsilon^{ll'}, \\
\overline{(\Sigma_R)^C} \Sigma_R &= \overline{(\Sigma_R)^C}_{ijkl} \Sigma_{Ri'j'k'l'} \epsilon^{ii'} \epsilon^{jj'} \epsilon^{kk'} \epsilon^{ll'}, \\
H^* H \Phi^* \Phi &= H^{*i} H_j \Phi^{*jkl} \Phi_{ikl}, \\
H^* H H \Phi &= H^{*i} H_j H_k \Phi_{ij'k'} \epsilon^{jj'} \epsilon^{kk'}, \\
H H \Phi \Phi &= H_i \Phi_{jkl} H_{i'} \Phi_{j'k'l'} \epsilon^{ij} \epsilon^{i'j'} \epsilon^{kk'} \epsilon^{ll'}, \\
H \Phi^* \Phi \Phi &= H_i \Phi^{*ijk} \Phi_{jlm} \Phi_{kl'm'} \epsilon^{ll'} \epsilon^{mm'}, \\
\Phi^* \Phi \Phi^* \Phi &= \Phi^{*ijk} \Phi_{i'jk} \Phi^{*i'j'k'} \Phi_{ij'k'}.
\end{aligned} \tag{4}$$

The light neutrino masses arise through tree-level dimension-nine operator and one-loop dimension-five operator. The diagrams displayed in Fig. 1 lead to two contributions to light neutrino masses

$$\begin{aligned}
(m_\nu)_{ij} &= (m_\nu)_{ij}^{\text{tree}} + (m_\nu)_{ij}^{\text{loop}} \\
&= \frac{-1}{6} (\lambda_4^*)^2 \frac{v_H^6}{\mu_\Phi^4} \sum_k \frac{Y_{ik} Y_{jk}}{M_k} + \frac{-5\lambda_5^* v_H^2}{24\pi^2} \sum_k \frac{Y_{ik} Y_{jk} M_k}{m_\Phi^2 - M_k^2} \\
&\quad \times \left[ 1 - \frac{M_k^2}{m_\Phi^2 - M_k^2} \ln \frac{m_\Phi^2}{M_k^2} \right],
\end{aligned} \tag{5}$$

determined by quartic couplings $\lambda_4$ and $\lambda_5$, respectively. Let us remind [19] that in the regime of comparable and light ($\sim 200$ GeV) masses for exotic fermions and scalars the tree contribution prevails, while the loop contribution dominates for heavy masses ($\sim 500$ GeV) and remains a sole contribution if additional discrete $Z_2$ symmetry forbids the $\lambda_4$ term in Eq. (3).

## 3. Scalar potential

After the electroweak symmetry breaking (EWSB) the neutral components of the scalar fields acquire a vacuum expectation value and read

$$H^0 = \frac{1}{\sqrt{2}} (v_H + h^0 + i\chi), \qquad \Phi^0 = \frac{1}{\sqrt{2}} (v_\Phi + \varphi^0 + i\eta). \tag{6}$$

If all couplings in the scalar potential are real the conditions for the minimum of the potential

$$\frac{\partial V_0(v_H, v_\Phi)}{\partial v_H} = 0, \qquad \frac{\partial V_0(v_H, v_\Phi)}{\partial v_\Phi} = 0 \tag{7}$$

give

$$\mu_H^2 = \lambda_1 v_H^2 + \left( \frac{1}{2}\lambda_2 + \frac{1}{6}\lambda_3 - \frac{2}{3}\lambda_5 \right) v_\Phi^2 + \frac{\sqrt{3}}{2} \lambda_4 v_H v_\Phi - \frac{\sqrt{3}}{9} \lambda_6 \frac{v_\Phi^3}{v_H}, \tag{8}$$

$$\mu_\Phi^2 = -\left( \frac{1}{2}\lambda_2 + \frac{1}{6}\lambda_3 - \frac{2}{3}\lambda_5 \right) v_H^2 - \frac{\sqrt{3}}{6} \lambda_4 \frac{v_H^3}{v_\Phi} + \frac{\sqrt{3}}{3} \lambda_6 v_H v_\Phi - \left( \lambda_7 + \frac{5}{9}\lambda_8 \right) v_\Phi^2. \tag{9}$$

The electroweak $\rho$ parameter is changed by $v_\Phi$ from the unit value to $\rho \simeq 1 + 6v_\Phi^2/v_H^2$ constraining the ratio $v_\Phi/v_H$ to be smaller then 0.015.

The mixing between singly charged and between neutral components of $H$ and $\Phi$ multiplets will occur after the EWSB. Let us illustrate this on the $h^0 - \varphi^0$ mass terms

$$\begin{aligned}
\mathcal{L}_{h^0\varphi^0} = &-\frac{1}{2} \left( -\mu_H^2 + 3\lambda_1 v_H^2 + \left( \frac{1}{2}\lambda_2 + \frac{1}{6}\lambda_3 - \frac{2}{3}\lambda_5 \right) v_\Phi^2 \right. \\
&\left. + \sqrt{3}\lambda_4 v_H v_\Phi \right) h^0 h^0 - \left( \left( \lambda_2 + \frac{1}{3}\lambda_3 - \frac{4}{3}\lambda_5 \right) v_H v_\Phi \right. \\
&\left. + \frac{\sqrt{3}}{2} \lambda_4 v_H^2 - \frac{1}{\sqrt{3}} \lambda_6 v_\Phi^2 \right) h^0 \varphi^0 \\
&- \frac{1}{2} \left( \mu_\Phi^2 + \left( \frac{1}{2}\lambda_2 + \frac{1}{6}\lambda_3 - \frac{2}{3}\lambda_5 \right) v_H^2 + \frac{2}{\sqrt{3}} \lambda_6 v_H v_\Phi \right. \\
&\left. + \left( 3\lambda_7 + \frac{5}{3}\lambda_8 \right) v_\Phi^2 \right) \varphi^0 \varphi^0.
\end{aligned} \tag{10}$$

The mass term that mixes $h^0$ and $\varphi^0$ is small because it is proportional either to $v_\Phi$ from $\lambda_{2,3,6}$ terms in Eq. (3), or to small lepton number violating couplings $\lambda_{4,5}$ dictated to be small from considerations of neutrino masses. Therefore we neglect the mixing between the components of $H$ and $\Phi$ multiplets and we neglect the terms of higher order in $v_\Phi/v_H$ whenever possible.

In the approximations given above, the masses of charged components of the quadruplet $\Phi$ are

$$\begin{aligned}
m^2(\Phi^+) &= \mu_\Phi^2 + \frac{1}{2}\lambda_2 v_H^2, \\
m^2(\Phi^-) &= \mu_\Phi^2 + \frac{1}{2}\lambda_2 v_H^2 + \frac{1}{3}\lambda_3 v_H^2, \\
m^2(\Phi^{--}) &= \mu_\Phi^2 + \frac{1}{2}\lambda_2 v_H^2 + \frac{1}{2}\lambda_3 v_H^2.
\end{aligned} \tag{11}$$

Their couplings to the Higgs boson relevant for the $h \to \gamma\gamma$ decay are

$$-\mathcal{L} = c_{\Phi^+} v_H h^0 \Phi^{+*} \Phi^+ + c_{\Phi^-} v_H h^0 \Phi^{-*} \Phi^- + c_{\Phi^{--}} v_H h^0 \Phi^{--*} \Phi^{--}, \tag{12}$$

where the newly introduced couplings

$$c_{\Phi^+} = \lambda_2, \qquad c_{\Phi^-} = \lambda_2 + \frac{2}{3}\lambda_3, \qquad c_{\Phi^{--}} = \lambda_2 + \lambda_3, \tag{13}$$

are expressed in terms of the quartic couplings $\lambda_2$ and $\lambda_3$ which, for simplicity, we assume to be equal in the following.

## 4. Higgs diphoton-decay width

In our model in addition to the dominant SM contributions from the $W$ boson and top quark loops to the $h \to \gamma\gamma$ decay rate only the charged scalars contribute substantially. Our exotic scalars are colorless so that the Higgs boson production through gluon