

Multi high charged scalars in the LHC searches and Majorana neutrino mass generations

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Abstract

One natural way to understand the excess of the measured $H \rightarrow \gamma\gamma$ rate over the standard model (SM) expectation at the Large Hadron Collider (LHC) is to have charged scalar bosons, existing in most of the SM extensions. Motivated by this LHC result, we explore if it also sheds light on solving the small neutrino mass generation problem. We concentrate on a class of models with high dimensional representations of scalars to realize Majorana neutrino masses at two-loop level without imposing any new symmetry. In these models, multi scalars with the electric charges higher than two are naturally expected, which not only enhance the $H \rightarrow \gamma\gamma$ rate, but provide more searching grounds at the LHC. In particular, the rate of $H \rightarrow Z\gamma$ also changes similar to that of the diphoton channel.

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Introduction— The Higgs doublet (H) not only breaks the electroweak gauge symmetries $SU(2)_L \times U(1)_Y$ to $U(1)_Q$ but also gives masses to the charged fermions in the standard model (SM). The main goals of the Large Hadron Collider (LHC) are the search of the Higgs boson (H) and the investigation of the electroweak symmetry breaking mechanism. Recently, a boson with its mass around 125 GeV has been discovered by both ATLAS [1] and CMS [2] collaborations, and its spin-parity property is further identified to be 0^+ based on the LHC full 2011+2012 dataset [3]. The scalar is most likely the Higgs particle as its properties are consistent with the SM predictions except the possible large production rate of $H \rightarrow \gamma\gamma$. In 2012, the excess of the rate in both experiments is around 2σ deviation from the SM prediction. Currently, the measured signal strength of Higgs to diphoton by the ATLAS collaboration is $\mu = 1.6_{-0.3}^{+0.3}$ [4], whereas the corresponding one by the CMS collaboration goes down to $\mu = 0.78_{-0.26}^{+0.28}$ [5]. In other words, the excess still survives in the ATLAS data but disappears in the CMS measurement. The elimination of the inconsistency between the two collaborations will rely on the future data accumulation and analysis in the LHC phase-II. In this study, we consider the possibility that some of the excess is sustained. Clearly, in such a case, it is still a call for new physics. One of natural mechanisms is to include new charged particles in the SM [6, 7], which would enhance the decay rate due to the new charged loop contributions.

On the other hand, neutrino oscillations observed by the solar, atmospheric, and reactor neutrino experiments reveal that neutrinos are massive but tiny and mix with each other, implying the need of new physics as there is no renormalizable Yukawa interaction for neutrinos related to H in the SM. Theoretical studies on model buildings to understand the small neutrino masses are enormously varied if it is allowed to freely invent new particles, scalars and/or fermions, beyond the SM ones. Without a new fermion, the only possibility of having non-zero neutrino masses is that neutrinos are Majorana fermions with the accidental global lepton symmetry in the SM either explicitly or spontaneously broken. Apart from the seesaw mechanism [8], models with tiny Majorana neutrino masses arising from quantum corrections have been proposed [9, 10]. It is clear that this type of theories is usually incorporated with some extension of the scalar sector in the SM. In particular, singly and doubly charged scalars are required in the radiative neutrino mass generation mechanisms at one and two-loop levels, respectively. As a result, the existence of new charged particles is a generic feature of the radiative neutrino models. These charged scalars in turn

would help for resolving the excess of the $H \rightarrow \gamma\gamma$ rate at the LHC as discussed in the literature [11]. However, in most of the above radiative models, since only $SU(2)_L$ singlet scalars are introduced, new physics effects are limited in the lepton sector, whereas those involving hadrons, such as the neutrinoless double beta decays believed as a benchmark of the Majorana nature of neutrinos, do not show up.

In Ref. [12], an unconventional neutrino mass generation model with a triplet scalar was proposed, leading to two doubly charged scalars with the neutrinoless double beta decays dominated by the short distance contribution. However, this type of models needs adding some *ad hoc* discrete symmetry to forbid the neutrino mass term at tree level [13], which makes these models less appealing. In this study, we examine the most general cases with arbitrary $SU(2)_L$ representations for scalars to avoid the tree level neutrino mass generation without imposing any new symmetry. It turns out that a class of models with high odd-dimensional scalars, which generate neutrino masses at two-loop level, exists naturally. Consequently, multi scalar bosons with the electric charges higher than 2-unit appear. These multi charged scalars clearly ensure us some rich phenomenologies at the LHC. We will explore them along with the lepton number violating processes.

Multi scalars— For all non-Higgs like scalars with non-trivial $SU(2)_L \times U(1)_Y$ quantum numbers [14], there are only three possible renormalizable Yukawa interactions, given by

$$f_{ab}\bar{L}_a^c L_b s, \quad y_{ab}\bar{\ell}_{R_a}^c \ell_{R_b} \Phi, \quad \text{and} \quad g_{ab}\bar{L}_a^c L_b T, \quad (1)$$

where L (ℓ_R) stands for the left-handed (right-handed) lepton, a or b denotes e, μ and τ , c represents the charged conjugation, s and Φ are $SU(2)_L$ singlet scalar fields with $Y = 2$ and $Y = 4$, and T is an $SU(2)_L$ triplet with the hypercharge $Y = 2$, respectively [15]. Note that the triplet scalar T can generate neutrino masses at tree level after developing its vacuum expectation value (VEV), known as the Type-II seesaw mechanism [16], while s and Φ are used in Zee [9] and Zee-Babu [10] models, respectively.

In Eq. (1), the third Yukawa interaction is the most troublesome as it generates neutrino masses at tree level with an extreme small value of the VEV or Yukawa couplings, which is obviously un-natural. Without introducing the triplet T , the interactions in Eq. (1) are precisely given by the Zee-Babu model [10], which has been extensively studied in the literature, in particular its phenomenology of the doubly charged scalar at the LHC. In this work, we consider a new class of models by adding a scalar field ξ with $Y = 2$ and a non-

trivial $SU(2)_L$ representation \mathbf{n} . To minimize our models, we disregard the singlet scalar s and keep the other singlet Φ so that only the second Yukawa interaction in Eq. (1) can exist at tree level.

The general scalar potential reads

$$\begin{aligned}
V(H, \xi, \Phi^{\pm\pm}) &= -\mu_H^2 |H|^2 + \lambda_H |H|^4 + \mu_\xi^2 |\xi|^2 + \lambda_\xi^\alpha |\xi|_\alpha^4 \\
&+ \mu_\Phi^2 |\Phi|^2 + \lambda_\Phi |\Phi|^4 + \lambda_{H\xi}^\beta (|H|^2 |\xi|^2)_\beta \\
&+ \lambda_{H\Phi} |H|^2 |\Phi|^2 + \lambda_{\xi\Phi} |\xi|^2 |\Phi|^2 + [\mu\xi\xi\Phi + \text{h.c.}], \tag{2}
\end{aligned}$$

where α, β are the short-handed notations denoting the possible invariant terms for higher representations in general. Also notice that all terms in the potential are self-hermitian except the last μ -term, which is related to the dynamics of the lepton number breaking to be discussed later. For the even dimensional representations, *i.e.* $\mathbf{n} = \mathbf{2}, \mathbf{4}, \mathbf{6} \dots$, the products $\xi\xi$ vanish since

$$\begin{aligned}
\xi\xi &= \epsilon_{ii'} \epsilon_{jj'} \epsilon_{kk'} \dots \xi_{ijk\dots} \xi_{i'j'k'\dots} \\
&= -\epsilon_{i'i} \epsilon_{j'j} \epsilon_{k'k} \dots \xi_{ijk\dots} \xi_{i'j'k'\dots} = 0, \tag{3}
\end{aligned}$$

due to the anti-symmetric matrix of ϵ_{ij} ($i, j = 1, 2$). Subsequently, we only need to consider the odd dimensional representations of ξ , *i.e.* $\mathbf{n} = \mathbf{3}, \mathbf{5}, \mathbf{7} \dots$. Since the triplet has been dropped out, the next minimal choice is $\mathbf{n} = \mathbf{5}$. From now on, we concentrate on this minimal one, *the quintuplet*, with $\xi = (\xi^{+++}, \xi^{++}, \xi^+, \xi^0, \xi^-)^T$. One shall bear in mind that the results can be easily extended to those with higher representations of ξ [17, 18]. In these cases, there are three and two irreducible terms for $|\xi|^4$ and $|H|^2 |\xi|^2$ respectively.

We now move to the lepton number (L) violation. In general, $U(1)_L$ can be either global or gauge symmetry. If the lepton number indeed comes from a global symmetry as that in the SM, the spontaneous symmetry breaking will generate a Nambu-Goldstone (NG) boson, usually called Majoron. In this model, the VEV of $v_\xi/\sqrt{2} = \langle \xi^0 \rangle$ breaks $U(1)_L$ spontaneously [19]. Since ξ is an $SU(2)$ multiplet, its corresponding Majoron has a direct coupling to the Z boson, which is strongly constrained by the LEP measurement of the invisible Z decay width. To resurrect it, we illustrate here by adding another $SU(2)_L$ triplet ($\mathbf{3}$) scalar field $\Delta = (\Delta^+, \Delta^0, \Delta^-)^T$ with $Y = 0$ to the model, which results in one additional non-hermitian term $\Delta\xi H^* H^*$. It is easy to see that the coexistences of $\bar{\ell}_R^c \ell_R \Phi$, $\mu\xi\xi\Phi$ and $\Delta\xi H^* H^*$ break the lepton number explicitly. Since v_ξ is constrained to be less than around

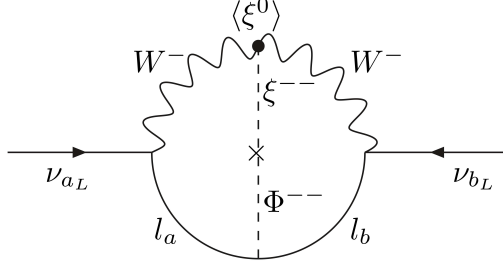


FIG. 1. Two-loop contributions to neutrino masses.

1 GeV by the ρ parameter [19, 20], it can be shown that the mixing angle between the neutral scalar components of ξ and H is of $\mathcal{O}(v_\xi/v)$. In practice, it is safe to ignore the mixing and treat the neutral component from the doublet field H as the SM-like Higgs. In addition, as ξ does not couple to the SM fermions directly, the single production of ξ^0 via the gluon-gluon fusion process as well as its W^+W^- fusion and W -associated productions are all suppressed by a factor of $(v_\xi/v)^2 \leq 10^{-4}$.

Neutrino masses— In the Zee-Babu model, neutrino masses are generated from the two-loop diagrams due to the couplings of $s^\pm s^\pm \Phi^\mp \mp$. whereas in our model without s^\pm , they can be induced from similar two-loop diagrams with $\Phi^{\pm\pm}$ coupling to $W^\mp W^\mp$ though the mixing of ξ as shown in Fig. 1. In this mechanism, neutrino masses are calculable, given by

$$m_{\nu_{ab}} \simeq \frac{g^4}{\sqrt{2}(4\pi)^4} m_a m_b v_\xi y_{ab} \sin 2\theta \times \left[\frac{1}{M_{P_1}^2} \log^2 \left(\frac{M_W^2}{M_{P_1}^2} \right) - \frac{1}{M_{P_2}^2} \log^2 \left(\frac{M_W^2}{M_{P_2}^2} \right) \right], \quad (4)$$

where $m_{a,b}$ correspond to charged lepton masses; P_1 and P_2 are the mass eigenstates of doubly charged scalars with θ representing their mixing angle and $M_{P_i} > M_W$ assumed. Note that the neutrino masses are suppressed by the two-loop factor, $SU(2)_L$ gauge coupling, charged lepton masses, mixing angle θ , and VEV of ξ , respectively, without fine-tuning Yukawa couplings y_{ab} . The model predicts the neutrino mass spectrum to be a normal hierarchy if one requires the perturbative bound on y_{ab} . Consequently, the neutrino mass matrix is given by

$$m_{\nu_{ab}} = U_{\text{PMNS}} m_{\nu_{diag}} U_{\text{PMNS}}^T \propto y_{ab}, \quad (5)$$

where $m_{\nu_{diag}} = \text{diag}(m_{\nu_1}, \sqrt{m_{\nu_1}^2 + \Delta m_{sol}^2}, \sqrt{m_{\nu_1}^2 + \Delta m_{atm}^2})$ with m_{ν_1} the lightest ν mass. From the neutrino oscillation data, one is able to pin down the neutrino parameters via the

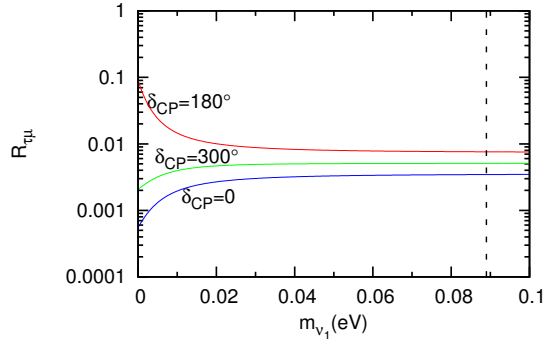


FIG. 2. $R_{\tau\mu}$ versus m_{ν_1} , where the blue, red and green curves correspond to $\delta_{CP} = 0, 180^\circ$ and 300° , respectively, with the global χ^2 analysis in Ref. [21], while the dashed black line represents the lower bound set by cosmology.

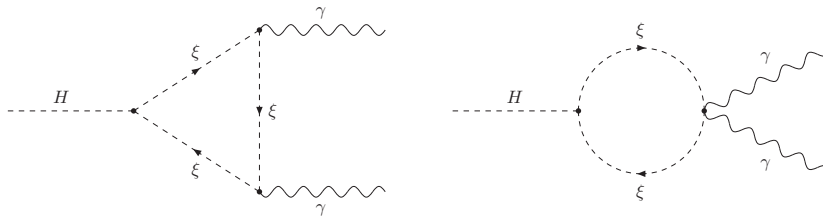


FIG. 3. Contributions to $H \rightarrow \gamma\gamma$ from charged scalar exchanges in the loops.

leptonic processes governed by y_{ab} . For example, the ratio $R_{\tau\mu} \equiv \frac{Br(\tau \rightarrow e\gamma)^*}{Br(\mu \rightarrow e\gamma)}$ [22] is related to the lightest neutrino mass, m_{ν_1} , as illustrated in Fig. 2, with the use of the latest neutrino oscillation data [21]. In principle, many variants of this quantity can be also defined from the leptonic rare decays or same-sign dilepton decays of $\Phi^{\pm\pm}$. Hence, our model provides a complementary way to determine neutrino parameters. Due to the indirect couplings between $\Phi^{\pm\pm}$ and $W^\mp W^\mp$, our model reveals a mechanism for possible large neutrinoless double β decays which do not directly involve the small Majorana neutrino masses unlike the Zee-Babu model.

LHC searches— In our minimal model with $\mathbf{n} = \mathbf{5}$, apart from one SM-like Higgs scalar with $m_H^2 = 2\lambda_H v^2$, there are one pseudo-scalar, three singly charged, two doubly charged, and one triply charged scalars with masses all around a few hundred GeV, which all satisfy the current experimental bounds [20]. Since ξ does not directly interact with the SM fermions, the Higgs production is not expected to be modified. However, as promised, the

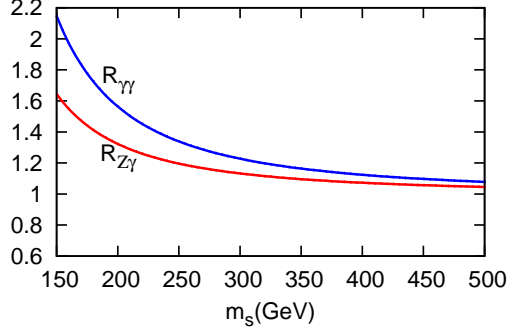


FIG. 4. $R_{\gamma\gamma} \equiv \Gamma(H \rightarrow \gamma\gamma)/\Gamma(H \rightarrow \gamma\gamma)_{\text{SM}}$ and $R_{Z\gamma} \equiv \Gamma(H \rightarrow Z\gamma)/\Gamma(H \rightarrow Z\gamma)_{\text{SM}}$ as functions of the degenerate mass factor m_s of the multi charged scalar states with $\mathbf{n} = \mathbf{5}$ and the universal trilinear coupling to Higgs, $\mu_s = -100$ GeV.

decay rate of $H \rightarrow \gamma\gamma$ receives extra contributions from the new charged scalars in the loops as shown in Fig. 3, given by

$$\begin{aligned} \Gamma(H \rightarrow \gamma\gamma) = & \frac{G_F \alpha^2 m_H^3}{128 \sqrt{2} \pi^3} \left| \sum_f N_f^c Q_f^2 A_{\frac{1}{2}}(\tau_f) + A_1(\tau_W) \right. \\ & \left. + \sum_{I_3} (I_3 + 1)^2 \frac{v}{2} \frac{\mu_s}{m_s^2} A_0(\tau_s) \right|^2, \end{aligned} \quad (6)$$

where the third term corresponds to those from the multi charged components of the dimension- \mathbf{n} scalar multiplet ξ with their electric charges $(I_3 + 1)$, I_3 runs from $-(\mathbf{n} - 1)/2$ to $(\mathbf{n} - 1)/2$, and μ_s is the trilinear coupling to the SM Higgs. The amplitudes $A_{0,\frac{1}{2},1}$ and the mass ratios $\tau_{f,W,s}$ are defined in Ref. [23]. Here, for simplicity, we have taken the same trilinear coupling μ_s and charged scalar mass m_s . The new contributions from the multi charged scalars interfere constructively with that of the SM if $\mu_s < 0$. As an illustration, we plot the ratio of $R_{\gamma\gamma} \equiv \Gamma(H \rightarrow \gamma\gamma)/\Gamma(H \rightarrow \gamma\gamma)_{\text{SM}}$ in Fig. 4 with a typical value of $\mu_s = -100$ GeV. It is clear that the excess for the diphoton decay rate reported by ATLAS can be easily explained by these multi charged scalars. Furthermore, the partial width of $H \rightarrow Z\gamma$ also receives new contributions from the multi charged scalar fields in our model. In Fig. 4, we plot $R_{Z\gamma} \equiv \Gamma(H \rightarrow Z\gamma)/\Gamma(H \rightarrow Z\gamma)_{\text{SM}}$ with $\mathbf{n} = \mathbf{5}$. Like the diphoton channel, $\Gamma(H \rightarrow Z\gamma)$ also gets enhanced [24] due to the new charged fields in comparison with the SM prediction. The simultaneous observations of $\gamma\gamma$ and $Z\gamma$ modes would help for discriminating the possible new physics beyond the SM [7].

The scalar multiplet components can be pair produced via $q\bar{q} \rightarrow \gamma/Z^* \rightarrow \xi\xi^*(\Phi^{++}\Phi^{--})$

m_ξ (GeV)	Signal [fb]	Signal [fb]
	$2\ell^+2\ell^- + \cancel{E}_T$	$\ell^+2\ell^-2j + \cancel{E}_T$
100	5.2×10^{-1}	3.3
200	8.2×10^{-2}	5.2×10^{-1}
300	2.1×10^{-2}	1.3×10^{-1}
400	6.7×10^{-3}	4.2×10^{-2}
500	2.5×10^{-3}	1.6×10^{-2}
	Backgrounds [fb]	Backgrounds [fb]
	$2W^+2W^- \quad W^+W^-Z$	$t\bar{t}W^- \quad 2W^-W^+2j \quad W^-Z2j$
	$2.2 \times 10^{-4} \quad 6.6 \times 10^{-2}$	$3.6 \times 10^{-1} \quad 6.7 \times 10^{-2} \quad 17.8$

TABLE I. Cross sections of the $\xi^{++}\xi^{--}$ -pair productions and the corresponding SM backgrounds with $\sqrt{s} = 8$ TeV, where $\ell = e$ or μ .

or $q\bar{q}' \rightarrow W^* \rightarrow \xi^{+++}\xi^{--}(\xi^{++}\xi^-, \xi^+\xi^0)$ at the LHC. Similar processes for the triplet Higgs productions have been investigated in the Type-II seesaw mechanism [25]. The $\xi^{++}\xi^{--}$ -pair production rate varies with m_ξ from 1 to 10^{-3} pb for $m_\xi = 100$ to 600 GeV with $\sqrt{s} = 8$ TeV. The lower mass limit of $\xi^{\pm\pm}$ has been set by CMS [26] and ATLAS [27] to be (382, 391, 395) and (409, 375, 398) GeV by assuming 100% branching ratio of ($e^\pm e^\pm$, $e^\pm \mu^\pm$, $\mu^\pm \mu^\pm$) final states, respectively. The bounds, of course, will reduce significantly if more leptonic modes are opened. One noticeable feature of our model is that there are two doubly charged states, $\xi^{\pm\pm}$ and $\Phi^{\pm\pm}$, which decay almost 100% into $W^\pm W^\pm$ and $l_R^\pm l_R^\pm$, respectively. In particular, the channels of $\xi^{\pm\pm} \rightarrow \xi^\pm \xi^\pm$ are kinematically forbidden. Therefore, the previous experimental bounds can be directly applied to $\Phi^{\pm\pm}$. For the $\xi^{++}\xi^{--}$ -pair productions, in Table I we list the cross sections of the signals with $m_\xi = 100 - 500$ GeV and the corresponding SM backgrounds, where we only consider the final lepton state to be e or μ due to the low tagging efficiency of τ . From the table, we see that there is only less than one signal event based on the current luminosity of about 20 fb^{-1} at the LHC if m_ξ is heavier than 200 GeV for four charged lepton modes. For the signals with more jets, we expect higher SM backgrounds although the kinematical cuts are usually used to enhance the signal significance. A detailed collider analysis is beyond the scope of this brief report. Furthermore, if we take $R_{\gamma\gamma} \simeq 1.2$ of the combined result from ATLAS and CMS

collaborations, the m_ξ is about 300 GeV.

The signals of the triply charged component in the quintuplet scalar are distinct and its production rate is about two times larger than the doubly charged one according to their coupling strengths. The triply charged scalar ξ^{+++} decays into three W bosons with the width

$$\Gamma(\xi^{+++} \rightarrow 3W^+) \simeq \frac{3g_W^6 v_\xi^2 m_\xi^5}{512\pi^3 m_W^6}, \quad (7)$$

which subsequently go into the final states with six fermions, where we have ignored the phase space suppression. Note that in the degenerate limit, all W bosons can be on-shell. In addition to the multiplet ξ , the doubly charged scalars $\Phi^{\pm\pm}$ also make the signals of our model more distinct at the LHC. If $m_{\xi^{+++}} > m_{\Phi^{++}}$, the channel $\xi^{+++} \rightarrow \Phi^{++} + W^{*+}$ opens up for the triply charged scalar through the mixings, which can be distinguished from the $3W^{*+}$ one as $\Phi^{\pm\pm}$ only decay to the same sign right-handed charged leptons with large center masses. By tagging the taus produced from Φ or W one can pin down the chirality of τ from its subsequent pion distribution spectrum [28]. Finally, we remark that the above results can be easily applied to other models with higher charged scalars.

Summary— We have proposed a common origin for the excess of the $H \rightarrow \gamma\gamma$ rate measured by the LHC and small neutrino masses based on models with some multi high charged scalars. We have shown that due to the charged scalars, both rates of $H \rightarrow \gamma\gamma$ and $H \rightarrow Z\gamma$ can be enhanced, while small Majorana neutrino masses are realized radiatively through the two-loop diagrams. Our models predict a normal hierarchy for the neutrino mass spectrum and allow a large rate for the neutrinoless double β decay. We have demonstrated that the multi high charged scalars provide many interesting searching grounds at the LHC.

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