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Doubly-charged scalar bosons from the doublet

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ABSTRACT

We consider the extended Higgs models, in which one of the isospin doublet scalar fields carries the hypercharge $Y = 3/2$. Such a doublet field $\Phi_{3/2}$ is composed of a doubly charged scalar boson as well as a singly charged one. We first discuss a simple model with $\Phi_{3/2}$ (Model I), and study its collider phenomenology at the LHC. We then consider a new model for radiatively generating neutrino masses with a dark matter candidate (Model II), in which $\Phi_{3/2}$ and an extra $Y = 1/2$ doublet as well as vector-like singlet fermions carry the odd quantum number for an unbroken discrete Z_2 symmetry. We also discuss the neutrino mass model (Model III), in which the exact Z_2 parity in Model II is softly broken. It is found that the doubly charged scalar bosons in these models show different phenomenological aspects from those which appear in models with a $Y = 2$ isospin singlet field or a $Y = 1$ triplet one. They could be clearly distinguished at the LHC.

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1. Introduction

Physics of electroweak symmetry breaking remains the last unknown part of the standard model (SM), and currently Higgs boson searches are underway at the Tevatron and the LHC. On the other hand, some phenomena beyond the SM such as neutrino oscillation and the existence of dark matter have been confirmed by experiments, and various models beyond the SM have been proposed to explain these phenomena. In such new physics models, non-minimal Higgs sectors are often introduced.

Charged scalar states generally appear in extended Higgs sectors. Although such charged states are in themselves new physics phenomena, by measuring their property the Higgs sector and thus the direction of new physics could be determined. In particular, doubly charged scalar states are a clear signature for Higgs sectors with non-standard representations such as isospin triplet fields with the hypercharge $Y = 1$ and singlet fields with $Y = 2$. Their phenomenological aspects strongly depend on the model, so that they can give important information to distinguish these models. The triplet fields are introduced in various models such as the left-right model [1], the littlest Higgs model [2] and the Type II seesaw model [3], while the singlet fields appear in various models for the grand unification and also in radiative seesaw models [4]. Phenomenology of these doubly charged states thus has been studied extensively.

There are, however, other representations which contain doubly charged scalar states. The simplest example is the isospin doublet scalar field $\Phi_{3/2}$ with $Y = 3/2$. Its phenomenology has hardly been studied. The field $\Phi_{3/2}$ may appear in a model for the coupling unification [5], or, as we shall discuss later in details, it can be used to build new versions of simple radiative seesaw models with a dark matter candidate [6–8] or without it [4,9].

In this Letter, we discuss phenomenology of $\Phi_{3/2}$ in renormalizable theories. Contrary to the $Y = 1$ triplet field Δ as well as the $Y = 2$ singlet S^{++} , the Yukawa coupling between $\Phi_{3/2}$ and charged leptons is protected by the chirality. In addition, the component fields of $\Phi_{3/2}$ are both charged and do not receive a vacuum expectation value (VEV) as long as electric charge is conserved. Hence, the field decays via the mixing with the other scalar representations which can decay into the SM particles or via some higher order couplings. This characteristic feature of $\Phi_{3/2}$ would give discriminative predictions at collider experiments. We therefore first study collider signatures of $\Phi_{3/2}$ at the LHC in the model (Model I) of an extension from the SM with an extra $Y = 1/2$ doublet and $\Phi_{3/2}$. We then present a new TeV scale model with $\Phi_{3/2}$, an extra $Y = 1/2$ doublet and vector-like singlet fermions (Model II), which contains a mechanism for radiatively generating tiny neutrino masses with a dark matter candidate under the exact Z_2 parity. We then briefly discuss phenomenology of the model under the current data from experiments for neutrino oscillation, lepton flavor violation (LFV), and dark matter. We also

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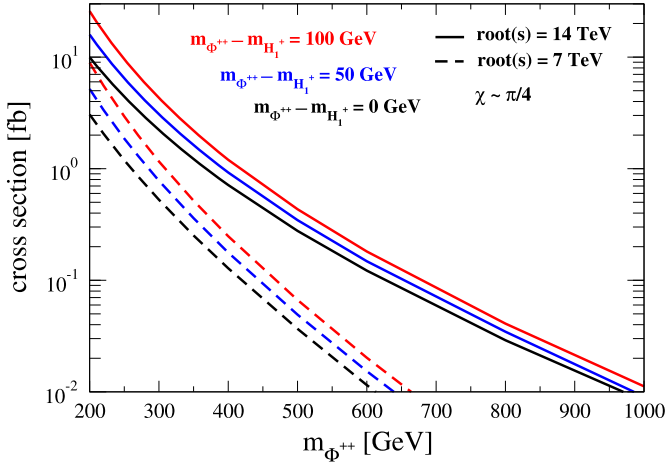


Fig. 1. Cross section of $pp \rightarrow XW^{++} \rightarrow X\Phi^{++}H_1^-$.

discuss the neutrino mass model (Model III), in which the Z_2 parity in Model II is softly broken.

2. Model I

The simplest model, where $\Phi_{3/2}$ is just added to the SM, can decay into SM particles only if lepton-number violating higher order operators are introduced [10]. Thus, we here consider the model in which $\Phi_{3/2}$ is added to the model with two $Y = 1/2$ Higgs doublet fields ϕ_1 and ϕ_2 (Model I). The singly charged scalar state in $\Phi_{3/2}$ can decay into the SM particles via the mixing with the physical charged state from the $Y = 1/2$ doublets. This model can be regarded as an effective theory of Model III which we discuss later, or it may be that of the model with an additional heavier Δ , in which the gauge coupling unification would be possible. In order to avoid flavor changing neutral current, a softly-broken Z_2 symmetry is imposed [11], under which the scalar fields are transformed as $\phi_1 \rightarrow \phi_1$, $\phi_2 \rightarrow -\phi_2$, and $\Phi_{3/2} \rightarrow -\Phi_{3/2}$.

The most general scalar potential is given by

$$\begin{aligned}
 V = & \sum_{i=1}^2 \mu_i^2 |\phi_i|^2 + (\mu_3^2 \phi_1^\dagger \phi_2 + \text{h.c.}) + \sum_{i=1}^2 \frac{1}{2} \lambda_i |\phi_i|^4 \\
 & + \lambda_3 |\phi_1|^2 |\phi_2|^2 + \lambda_4 |\phi_1^\dagger \phi_2|^2 + \frac{1}{2} [\lambda_5 (\phi_1^\dagger \phi_2)^2 + \text{h.c.}] \\
 & + \mu_\Phi^2 |\Phi_{3/2}|^2 + \frac{1}{2} \lambda_\Phi |\Phi_{3/2}|^4 + \sum_{i=1}^2 \rho_i |\phi_i|^2 |\Phi_{3/2}|^2 \\
 & + \sum_{i=1}^2 \sigma_i |\phi_i^\dagger \Phi_{3/2}|^2 + [\kappa (\Phi_{3/2}^\dagger \phi_1) (\phi_2 \cdot \phi_1) + \text{h.c.}], \quad (1)
 \end{aligned}$$

where the Z_2 symmetry is softly broken at the μ_3^2 term. We neglect the CP violating phase for simplicity. The scalar doublets ϕ_1 , ϕ_2 and $\Phi_{3/2}$ are parameterized as

$$\phi_i = \begin{bmatrix} w_i^+ \\ \frac{1}{\sqrt{2}}(h_i + v_i + iz_i) \end{bmatrix} \quad (i = 1, 2), \quad \Phi_{3/2} = \begin{bmatrix} \Phi^{++} \\ \Phi^+ \end{bmatrix},$$

where the VEVs v_i satisfy $\sqrt{v_1^2 + v_2^2} = v \simeq 246$ GeV. Mass matrices for the neutral components are diagonalized as in the same way as those in the usual two Higgs doublet model (2HDM) with ϕ_1 and ϕ_2 . The mass eigenstates h and H for CP-even states are obtained by diagonalizing the mass matrix by the angle α . By the angle β ($\tan \beta \equiv v_2/v_1$), the mass eigenstates for the CP-odd states z and

A are obtained, where z is the Nambu–Goldstone (NG) boson and A is the CP-odd Higgs boson. For simplicity $\sin(\beta - \alpha) = 1$ is taken such that h is the SM-like Higgs boson [12]. The existence of $\Phi_{3/2}$ affects the singly charged scalar sector. The mass eigenstates are obtained by mixing angles β and χ as

$$\begin{bmatrix} w^\pm \\ H_1^\pm \\ H_2^\pm \end{bmatrix} = \begin{bmatrix} 1 & 0 & 0 \\ 0 & c_\chi & s_\chi \\ 0 & -s_\chi & c_\chi \end{bmatrix} \begin{bmatrix} c_\beta & s_\beta & 0 \\ -s_\beta & c_\beta & 0 \\ 0 & 0 & 1 \end{bmatrix} \begin{bmatrix} w_1^\pm \\ w_2^\pm \\ \Phi^\pm \end{bmatrix}, \quad (2)$$

where $c_\theta = \cos \theta$ and $s_\theta = \sin \theta$, w^\pm are the NG bosons absorbed by the longitudinal component of the W^\pm bosons. H_1^\pm and H_2^\pm are physical mass eigenstates with the masses $m_{H_1^\pm}$ and $m_{H_2^\pm}$.

The Yukawa couplings for charged states are given by

$$\begin{aligned}
 \mathcal{L}_Y = & -\frac{\sqrt{2}V_{KM}^{ij}}{v} \bar{u}^i (m_{ui} \xi_A^u d_L^j + m_{dj} \xi_A^d d_R^j) \phi^+ \\
 & -\frac{\sqrt{2}m_{\ell i} \xi_A^\ell}{v} \bar{\nu}^i \ell_R^i \phi^+ + \text{h.c.} \quad (\phi^+ = H_1^+ \cos \chi - H_2^+ \sin \chi),
 \end{aligned}$$

where the coupling parameters $\xi_A^{u,d,\ell}$ [13] depend on the Z_2 charges of quarks and leptons [14]. We are interested in the light charged scalar bosons such as $\mathcal{O}(100)$ GeV. To satisfy the $b \rightarrow s\gamma$ data [15], we choose the Type-I Yukawa interaction with $\tan \beta \gtrsim 2$. Assuming $m_{H_{1,2}^\pm} < m_t + m_b$ and $m_{H_2^\pm} - m_{H_1^\pm} < m_Z$, the branching ratios for the main decay modes are evaluated as $B(H_{1,2}^\pm \rightarrow \tau^\pm \nu) \sim 0.7$ and $B(H_{1,2}^\pm \rightarrow cs) \sim 0.3$ when $\chi \neq 0$.

At the LHC, Φ^{++} can be tested by using various processes such as the pair production and the associated production with H_1^- or H_2^- . We here discuss an interesting signal via the process $u\bar{d} \rightarrow W^{++} \rightarrow \Phi^{++}H_{1,2}^-$. The cross section is shown in Fig. 1. We may examine this process, for example, by the decay $\Phi^{++} \rightarrow H_{1,2}^+ W^+ \rightarrow \tau^+ \ell^+ \nu$ with $H_{1,2}^- \rightarrow jj$, when $m_{\Phi^{++}} > m_{H_1^\pm} + m_W$ with $m_{H_{1,2}^\pm} < m_t + m_b$ and $m_{H_2^\pm} - m_{H_1^\pm} < m_Z$, where $m_{\Phi^{++}}$ is the mass of Φ^{++} . The signal is then $\tau^+ \ell^+ jj$ plus a missing transverse momentum \cancel{E}_T ($\ell^+ = e^+$ or μ^+). The signal cross section for $\tau^\pm \ell^\pm jj \cancel{E}_T$ is evaluated as 4.0 fb (1.3 fb) for $\sqrt{s} = 14$ TeV (7 TeV) for $m_{H_1^\pm} = 100$ GeV, $m_{H_2^\pm} = 150$ GeV, $m_{\Phi^{++}} = 200$ GeV and $\chi \simeq \pi/4$.

The mass for Φ^{++} can be determined from the Jacobian peak [16] in the distribution of the transverse mass, $M_T(\tau^+ \ell^+ \cancel{E}_T) = \sqrt{2p_T^{\tau\ell} \cancel{E}_T (1 - \cos \varphi)}$, where φ is the azimuthal angle between the transverse momentum $p_T^{\tau\ell}$ of the dilepton system and \cancel{E}_T . We show numerical results for the scenario with $m_{H_1^\pm} = 100$ GeV, $m_{H_2^\pm} = 150$ GeV, $m_{\Phi^{++}} = 200$ GeV, $\chi \simeq \pi/4$, $\tan \beta = 3$, $\sin(\beta - \alpha) = 1$ and $m_H = m_A = 127$ GeV, where m_H and m_A represent the masses of H and A , respectively. The potential is then approximately custodial symmetric, so that the rho parameter constraint is satisfied with the mass of the SM-like Higgs boson h to be 120 GeV. The end point in Fig. 2 (left) indicates $m_{\Phi^{++}}$, where the event number is taken to be 1000. One might think that the final decay products from the τ lepton should be discussed. We stress that the endpoint at $m_{\Phi^{++}}$ also appears in the distribution of $M_T(\ell^+ \ell^+ \cancel{E}_T)$ obtained from the leptonic decay of the τ^+ . The cross section for the $\ell^+ \ell^+ jj \cancel{E}_T$ signal is about 1.3 fb for $\sqrt{s} = 14$ TeV (0.45 fb for $\sqrt{s} = 7$ TeV). Furthermore, masses of singly charged Higgs bosons can also be measured by the distribution of $M_T(jj)$. In Fig. 2 (right), the two Jacobian peaks at 100 and 150 GeV correspond to $m_{H_1^\pm}$ and $m_{H_2^\pm}$, respectively, where the event number is taken to be 1000. The SM background for $\ell^+ \ell^+ jj \cancel{E}_T$, which mainly comes from $u\bar{d} \rightarrow W^+ W^+ jj$, is 3.95 fb for $\sqrt{s} = 14$ TeV (0.99 fb for $\sqrt{s} = 7$ TeV). The cross section of the

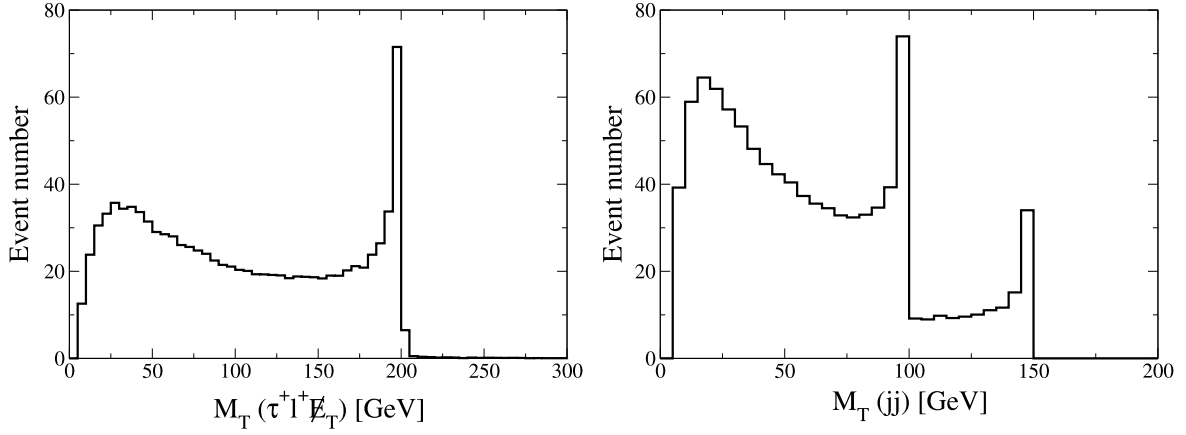


Fig. 2. (Left) The transverse mass distribution for the $\tau^+\ell^+\cancel{E}_T$ system for the signal. (Right) That for the jj system. The event number is taken to be 1000 for both figures.

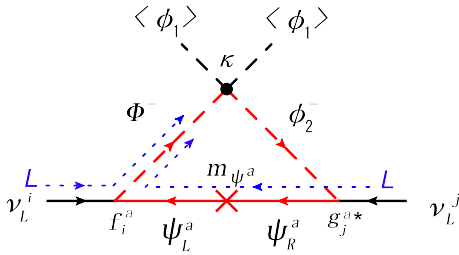


Fig. 3. Neutrino mass diagram.

background is comparable to that for the signal before kinematic cuts. There is no specific kinematical structure in the $\ell^+\ell^+$ distribution in the background. All the charged scalar states can be measured simultaneously via this process unless their masses are too heavy if sufficient number of the signal event remains after kinematic cuts. While the detection at the LHC with 300 fb^{-1} may be challenging, it could be much better at the upgraded version of the LHC with 3000 fb^{-1} .

3. Model II

We here present a new model in which $\Phi_{3/2}$ is introduced to naturally generate tiny neutrino masses at one-loop level. To this end, we again consider the scalar sector with ϕ_1 , ϕ_2 and $\Phi_{3/2}$. In addition, we introduce two isospin singlet Dirac fermions ψ^a ($a = 1, 2$) with $Y = -1$. We impose the exact (unbroken) Z_2 parity, under which ϕ_2 , $\Phi_{3/2}$ and ψ^a are odd while all the SM particles including ϕ_1 are even. This Z_2 parity plays a role to forbid mixing terms of $\bar{\ell}_R \psi_L^a$ as well as couplings of $\bar{L}_L \phi_1 \psi_R^a$ and $\bar{L}_L \phi_2 \ell_R$, and to guarantee the stability of a dark matter candidate; i.e., the lightest neutral Z_2 odd particle. Lepton numbers $L = -2$ and $+1$ are respectively assigned to $\Phi_{3/2}$ and ψ^a .

The scalar potential coincides that in Eq. (1) but $\mu_3^2 = 0$ due to the exact Z_2 parity. Without $\Phi_{3/2}$, the scalar sector is that of the inert doublet model [17], in which only ϕ_1 receives the VEV yielding the SM-like Higgs boson h , while ϕ_2 gives Z_2 -odd scalar bosons H , A and H^\pm . Including $\Phi_{3/2}$, H^\pm can mix with Φ^\pm diagonalized by the angle χ in Eq. (2) with $\beta = 0$. Masses and interactions for ψ^a are given by

$$\mathcal{L}_Y = m_{\psi^a} \bar{\psi}_L^a \psi_R^a + f_i^a (\bar{L}_L^i)^c \cdot \Phi_{3/2} \psi_L^a + g_i^a \bar{L}_L^i \phi_2 \psi_R^a + \text{h.c.} \quad (3)$$

The neutrino masses are generated via the one-loop diagram in Fig. 3. The flow of the lepton number is also indicated in the figure. The source of lepton number violation (LNV) is the coupling κ .

This is similar to the model by Zee [9], although the diagram looks similar to the model by Ma [7] where Majorana masses of right-handed neutrinos ν_R is the origin of LNV. For $m_{\psi^a} \gg m_{H_1^\pm}, m_{H_2^\pm}$, the mass matrix can be calculated as

$$(\mathcal{M}_\nu)_{ij} \simeq \sum_{a=1}^2 \frac{1}{16\pi^2} \frac{1}{2m_{\psi^a}} (f_i^a g_j^{a*} + f_j^a g_i^{a*}) \frac{v^2 \kappa}{m_{H_2^\pm}^2 - m_{H_1^\pm}^2} \times \left(m_{H_2^\pm}^2 \log \frac{m_{\psi^a}^2}{m_{H_2^\pm}^2} - m_{H_1^\pm}^2 \log \frac{m_{\psi^a}^2}{m_{H_1^\pm}^2} \right).$$

For $m_\psi \sim 1 \text{ TeV}$, $m_{H_1^\pm} \sim m_{H_2^\pm} \sim \mathcal{O}(100) \text{ GeV}$, and $f_i^a \sim g_i^a \sim \kappa \sim \mathcal{O}(10^{-3})$, the scale of neutrino masses ($\sim 0.1 \text{ eV}$) can be generated. The bound from LFV processes such as $\mu \rightarrow e\gamma$ [18] can easily be satisfied. The neutrino data can be reproduced by introducing at least two fermions ψ^1 and ψ^2 . The lightest Z_2 odd neutral Higgs boson (either H or A) is a dark matter candidate [17]. Assuming that H is the lightest, its thermal relic abundance can explain the WMAP data [19] by the s -channel process $HH \rightarrow h \rightarrow b\bar{b}$ (or $\tau^+\tau^-$). The t -channel process $HH \rightarrow \bar{\ell}_L \ell_L$ with ψ_R mediation is negligible. The direct search results can also be satisfied.

Finally, we comment on the collider signature in Model II. $\Phi_{3/2}$ is Z_2 odd, so that its decay product includes the dark matter H . For $m_H = 50 \text{ GeV}$, the mass of h would be about 115 GeV to satisfy the WMAP data [19]. We then consider the parameter set; $m_{\Phi^{++}} = 230 \text{ GeV}$, $m_{H_2^\pm} = 150 \text{ GeV}$, $m_{H_1^\pm} \simeq m_A = 149 \text{ GeV}$ and $\chi = 0.1$ to satisfy the neutrino data and the LFV data. The signal at the LHC would be $W^+W^+W^-\cancel{E}_T$ via $u\bar{d} \rightarrow \Phi^{++}H_1^- \rightarrow (H_1^+W^+)(W^-H) \rightarrow (HW^+W^+)(W^-H)$. The cross section of $W^\pm W^\pm W^\mp \cancel{E}_T$ is 23 fb for $\sqrt{s} = 14 \text{ TeV}$ (7.3 fb for $\sqrt{s} = 7 \text{ TeV}$). The main background comes from $W^\pm W^\pm W^\mp$, and the cross section is 135 fb for $\sqrt{s} = 14 \text{ TeV}$ (76 fb for $\sqrt{s} = 7 \text{ TeV}$). The signal background ratio is not too small at all, and we can expect the signal would be detected after appropriate kinematic cuts.

4. Discussions and conclusions

Let us discuss the exact (unbroken) Z_2 parity in Model II. In radiative seesaw models with ν_R [6–8], the Majorana masses are the source of LNV, and an exact Z_2 has to be imposed to forbid the neutrino Yukawa coupling. On the contrary, in Model II, the source of LNV is the κ term in Eq. (1), so that the canonical seesaw mechanism [20] cannot occur even without such an exact Z_2 parity. Thus, in Model II the exact Z_2 parity is not necessarily important for radiative generation of neutrino masses. Therefore, we

may consider another model (Model III), in which the exact Z_2 parity is softly broken in Model II. Then, in addition to the terms in Eq. (3), new terms of $\bar{\psi}_L^a \ell_R^i$ appear. They cause LFV processes such as $\mu \rightarrow eee$ or $\mu \rightarrow e\gamma$. By setting m_{ψ^a} to be at TeV scales with smaller mixing parameters $\tilde{m}_{\psi\ell}^{ai}$ of $\bar{\psi}_L^a \ell_R^i$, such LFV processes can be suppressed to satisfy the current data. Note that the tree level LFV process such as $\mu \rightarrow eee$ via $\bar{\psi}_L - \ell_R$ mixing is multiplied by a suppression factor like $m_\ell^2 \tilde{m}_{\psi\ell}^2 / (v^2 m_\psi^2)$. Neutrino masses are generated by not only the diagram in Fig. 3 but also additional diagrams where ℓ_R^i instead of ψ_R^a and ϕ^- instead of ϕ_2^- are in the loop. Phenomenology of the Higgs sector in Model III coincides that in Model I, so that the same signal shown in Fig. 2 can be used to identify the $\Phi_{3/2}$ field.

We have briefly discussed collider phenomenology of $\Phi_{3/2}$ in Models I–III. We comment on the difference from models with S^{++} or $\Delta = (\Delta^{++}, \Delta^+, \Delta^0)$. First, the production process $u\bar{d} \rightarrow W^{++} \rightarrow \varphi^{++}\varphi^-$ (φ : a scalar component) can be useful to test $\Phi_{3/2}$ and Δ [21], while S^{++} does not contribute to this process because of no weak gauge coupling. In the model with a Δ , although the signal strongly depends on the parameters [22], $\ell^+\ell^+jj\cancel{E}_T$ can be important for $m_{\Delta^{++}} - m_{\Delta^+} > \mathcal{O}(1)$ GeV when the hadronic decay mode of the singly charged state is substantial. In this case, masses of doubly and singly charged scalars can be measured from the end points in $M_T(\ell^+\ell^+\cancel{E}_T)$ and $M_T(jj)$ distributions, similarly to Models I and III. Otherwise, a peak in the invariant mass distribution $M(\ell^+\ell^+)$ would be seen at $m_{\Delta^{++}}$ in the triplet models when the VEV of the triplet field is sufficiently small [22], while there is no such a peak in the model with $\Phi_{3/2}$. Second, the production mechanism $u\bar{d} \rightarrow W^{*+} \rightarrow \varphi^+\varphi^0$ can be useful to study multi-doublet models [23] and the model with Δ as well as that with $\Phi_{3/2}$ with an extra doublet. There is no such process in the model with S^\pm . The signal strongly depends on the model; i.e., $bb\tau\cancel{E}_T$ for the Type-II 2HDM [23], $\tau\tau\tau\cancel{E}_T$ for the Type-X 2HDM [13], and some others in the model with Δ . Therefore, extracting properties of doubly and singly charged Higgs bosons by using these processes as well as other various specific processes, we can discriminate the Higgs sectors at the LHC (and its luminosity-upgraded version) to a considerable extent.

To summarize, we have studied various aspects of $\Phi_{3/2}$ including the signature at the LHC in a few models. New TeV-scale models with $\Phi_{3/2}$ have been presented for generating tiny neutrino masses, one of which also contains dark matter candidates. We have found that $\Phi_{3/2}$ in these models shows discriminative and testable aspects at the LHC and its luminosity-upgraded version, so that models with $\Phi_{3/2}$ would be distinguishable from the other models with doubly charged scalar states. The details are discussed elsewhere [24].

Note added

After this Letter was completed, a paper [25] appeared, where same sign dilepton resonances were discussed for Φ^{++} together with Δ^{++} and S^{++} .

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