ЖЭТФ, 2015, том 147, вып. 3, стр. 426–432

DOUBLE HIGGS PRODUCTION AT LHC, SEE-SAW TYPE-II AND GEORGI–MACHACEK MODEL

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Received September 25, 2014

The double Higgs production in the models with isospin-triplet scalars is studied. It is shown that in the see-saw type-II model, the mode with an intermediate heavy scalar, $pp \to H + X \to 2h + X$, may have the cross section that is comparable with that in the Standard Model. In the Georgi–Machacek model, this cross section could be much larger than in the Standard Model because the vacuum expectation value of the triplet can be large.

Contribution for the JETP special issue in honor of V. A. Rubakov's 60th birthday

DOI: 10.7868/S0044451015030040

1. INTRODUCTION

After the discovery of the Higgs-BE boson at the LHC [1], the next steps to check the Standard Model (SM) are the measurement of the coupling constants of the Higgs boson with other SM particles ($t\bar{t}, WW, ZZ, b\bar{b}, \tau\bar{\tau}, \ldots$) with better accuracy and the measurement of the Higgs self-coupling that determines the shape of the Higgs potential. In the SM, the triple and quartic Higgs couplings are predicted in terms of the known Higgs mass and vacuum expectation value. Deviations from these predictions would mean the existence of a New Physics in the Higgs potential. The triple Higgs coupling can be measured at the LHC in double Higgs production, in which the gluon fusion dominates: $gg \to hh$. However, the $2h$ production cross section is very small. According to [2], the cross section at $\sqrt{s} = 14$ TeV is $\sigma^{NNLO}(gg \to hh) = 40.2$ fb with a 10–15 % accuracy. For the final states with reasonable signal/background ratios (such as $hh \to b\bar{b}\gamma\gamma$), only at the HL-LHC with the integrated luminosity $\int \mathcal{L}\, dt = 3000\ \text{fb}^{-1}$ will the double Higgs production be found and the triple Higgs coupling will be measured [3][1)]. We seek the extensions of the SM Higgs sector in which the double Higgs production is enhanced.

One of the well-motivated examples of a nonminimal Higgs sector is provided by the see-saw type-II mechanism of neutrino mass generation [7]. In this mechanism, a scalar isotriplet $(\Delta^{++}, \Delta^{+}, \Delta^{0})$ with the hypercharge $Y_\Delta = 2$ is added to the SM. The vacuum expectation value (vev) of the neutral component v_Δ generates Majorana masses of the left-handed neutrinos. There are two neutral scalar bosons in the model: the light one, in which the doublet Higgs component dominates and which should be identified with the particle discovered at the LHC ($h; M_h = 125$ GeV), and the heavy one, in which the triplet Higgs component dominates (H). The neutrino masses equal $f_i v_\Delta$, where f_i ($i = 1, 2, 3$) originates from Yukawa couplings of the

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1) The decays into $b\bar{b}\tau\bar{\tau}$ and $b\bar{b}W^+W^-$ final states can be even more promising for measuring the triple Higgs coupling [4, 5] (see also [6]).

Higgs triplet with the lepton doublets. If neutrinos are light due to a small value of v_Δ while f_i are of the order of unity, then H decays into neutrino pairs. Three states, $H^{\pm\pm}$ (or $\Delta^{\pm\pm}$), $H^{\pm}$, and H, are almost degenerate in the model considered in Sec. 2, and the absence of the same-sign dileptons at the LHC from $H^{\pm\pm} \to l^{\pm}l^{\pm}$ decays provides the lower bound $m_H > 400$ GeV [8]. We are interested in the opposite case where v_Δ reaches the maximum allowed value while neutrinos are light because of small values of f_i. In this case, $H \to hh$ can be the dominant decay mode of a heavy neutral Higgs boson. In this way, we obtain an additional mechanism of the double h production at the LHC.

The bound $m_{H^{++}} > 400$ GeV [8] cannot be applied now because $H^{\pm\pm}$ mainly decays into the same-sign diboson [9]. We only need H to be heavy enough for the $H \to hh$ decay to occur. This case is analyzed in Sec. 2. The invariant mass of the additionally produced hh state peaks at $(p_1+p_2)^2 = m_H^2$, which is a distinctive feature of the proposed mechanism (see also [10, 11]).

H contains a small admixture of the isodoublet state, which makes gluon fusion a dominant mechanism of H production at the LHC. The admixture of the isodoublet component in H equals approximately $2v_\Delta/v$, where $v \approx 250$ GeV is the vacuum expectation value of the neutral component of the isodoublet, and in Sec. 2, for $\sqrt{s} = 14$ TeV and $M_H = 300$ GeV, we obtain $\sigma(gg \to H) \approx 25$ fb. Taking into account that the branching ratio $H \to hh$ is about 80 %, we obtain a 50 % enhancement of double Higgs production in comparison with the SM.

Since the nonzero value of v_Δ violates the well-checked equality of the strengths of charged and neutral currents at the tree level,

$$\frac{g^2/M_W^2}{\bar{g}^2/M_Z^2} = 1 + 2\frac{v_\Delta^2}{v^2}, \tag{1}$$

v_Δ should be less than 5 GeV (see Sec. 2). The $gg \to H$ cross section was estimated numerically for the maximum allowed value $v_\Delta = 5$ GeV when the isodoublet admixture is about 5 %.

The bound $v_\Delta < 5$ GeV is removed in the Georgi–Machacek (GM) model [12], where, in addition to $\boldsymbol{\Delta}$, a scalar isotriplet with $Y = 0$ is introduced. If the vev of the neutral component of this additional field equals v_Δ, then we have just one in the r.h.s. of Eq. (1): the correction proportional to v_Δ^2 is canceled. Thus, v_Δ can be much larger than 5 GeV. The bounds on v_Δ come from the measurement of the 125 GeV Higgs boson couplings to vector bosons and fermions, which would deviate from their SM values: $c_i \to c_i\left[1 + a_i\,(v_\Delta/v)^2\right]$.

The consideration of an enhancement of $2h$ production in the GM variant of the see-saw type-II model is presented in Sec. 3. Because the current accuracy of the measurement of c_i values in h production and decay is poor, v_Δ as large as 50 GeV is allowed and $\sigma(gg \to H)$ can reach 2 pb, which makes it accessible with the integrated luminosity $\int \mathcal{L}dt = 300$ fb^{-1} prior to the HL-LHC run. We summarize our results in the Conclusions.

2. DOUBLE *h* PRODUCTION IN *H* DECAYS AT THE LHC

2.1. Scalar sector of the see-saw type-II model

In this subsection, we present the necessary formulas (see [13] for a detailed description). In addition to the SM isodoublet field

$$\Phi \equiv \begin{bmatrix}\Phi^+\\ \Phi^0\end{bmatrix} \equiv \begin{bmatrix}\Phi^+\\ \frac{1}{\sqrt{2}}(v+\varphi+i\chi)\end{bmatrix}, \tag{2}$$

an isotriplet is introduced in the see-saw type-II model:

$$\Delta \equiv \frac{\boldsymbol{\Delta}\cdot\boldsymbol{\sigma}}{\sqrt{2}} = \begin{bmatrix}\Delta^3/\sqrt{2} & (\Delta^1 - i\Delta^2)/\sqrt{2}\\ (\Delta^1 + i\Delta^2)/\sqrt{2} & -\Delta^3/\sqrt{2}\end{bmatrix} \equiv \begin{bmatrix}\delta^+/\sqrt{2} & \delta^{++}\\ \delta^0 & -\delta^+/\sqrt{2}\end{bmatrix}, \quad \delta^0 = \frac{1}{\sqrt{2}}(v_\Delta+\delta+i\eta). \tag{3}$$

Here, $\boldsymbol{\sigma}$ are the Pauli matrices.

The scalar sector kinetic terms are

$$\mathcal{L}_{\text{kinetic}} = |D_\mu\Phi|^2 + \mathrm{Tr}\left[(D_\mu\Delta)^\dagger(D_\mu\Delta)\right], \tag{4}$$

where

$$D_\mu\Phi = \partial_\mu\Phi - i\frac{g}{2}A_\mu^a\sigma^a\Phi - i\frac{g'}{2}B_\mu\Phi, \tag{5}$$

$$D_\mu\Delta = \left[\partial_\mu\Delta^a + g\varepsilon^{abc}A_\mu^b\Delta^c - ig'B_\mu\Delta^a\right]\frac{\sigma^a}{\sqrt{2}} = \partial_\mu\Delta - i\frac{g}{2}\left[A_\mu^a\sigma^a, \Delta\right] - ig'B_\mu\Delta. \tag{6}$$

The hypercharge $Y_\Phi = 1$ was substituted for the isodoublet and $Y_\Delta = 2$ for the isotriplet. The terms quadratic in the vector boson fields are as follows:

$$\mathcal{L}_{V^2} = g^2 \left|\delta^0\right|^2 W^+W^- + \frac{1}{2}g^2 \left|\Phi^0\right|^2 W^+W^- + \\ + \bar{g}^2 \left|\delta^0\right|^2 Z^2 + \frac{1}{4}\bar{g}^2 \left|\Phi^0\right|^2 Z^2. \quad (7)$$

The vector boson masses are

$$\begin{aligned} M_W^2 &= \frac{g^2}{4}\left(v^2 + 2v_\Delta^2\right), \\ M_Z^2 &= \frac{\bar{g}^2}{4}\left(v^2 + 4v_\Delta^2\right). \end{aligned} \quad (8)$$

For the ratio of vector boson masses, neglecting the radiative corrections from the isotriplet (not a bad approximation as far as the heavy triplet decouples), we obtain

$$\frac{M_W}{M_Z} \approx \left(\frac{M_W}{M_Z}\right)_{\rm SM}\left(1 - \frac{v_\Delta^2}{v^2}\right). \quad (9)$$

Comparing the result of the SM fit [14, p. 145] $M_W^{\rm SM} = 80.381$ GeV with the experimental value $M_W^{\rm exp} = 80.385(15)$ GeV, we obtain the following upper bound at the 3σ level:

$$v_\Delta < 5 \text{ GeV}. \quad (10)$$

Because the cross sections we are interested in are proportional to v_Δ^2, we use the upper bound $v_\Delta = 5$ GeV for numerical estimates in this section.

From the numerical value of the Fermi coupling constant in muon decay, we obtain

$$v^2 + 2v_\Delta^2 = (246 \text{ GeV})^2, \quad (11)$$

and hence for $v_\Delta \lesssim 5$ GeV, the value $v = 246$ GeV can be safely used in deriving (10).

The scalar potential has the form

$$V(\Phi,\Delta) = -\frac{1}{2}m_\Phi^2\left(\Phi^\dagger\Phi\right) + \frac{\lambda}{2}\left(\Phi^\dagger\Phi\right)^2 + \\ + M_\Delta^2 \mathrm{Tr}\left[\Delta^\dagger\Delta\right] + \frac{\mu}{\sqrt{2}}\left(\Phi^T i\sigma^2\Delta^\dagger\Phi + \text{H.c.}\right), \quad (12)$$

which is a truncated version of the most general renormalizable potential (see, e. g., [15, Eq. (2.6)]). We may simply suppose that the coupling constants that multiply the omitted terms in the potential ($\lambda_1, \lambda_2, \lambda_4$, and λ_5) are small[2)]. In the case of the SM, only the first line in (12) remains; the mass of the Higgs boson equals $m_\Phi = 125$ GeV while its expectation value is $v^2 \approx \approx m_\Phi^2/\lambda \approx (246 \text{ GeV})^2$, $\lambda \approx 0.25$.

[2)] We note that a relatively large value of λ_5 leads to considerable splitting of the masses of triplet states. If the cascade decay $H \to H^+W^-$ becomes allowed, it greatly diminishes the $H \to hh$ branching ratio [16].

At the minimum of (12), the equations

$$\begin{aligned} \frac{1}{2}m_\Phi^2 &= \frac{1}{2}\lambda v^2 - \mu v_\Delta, \\ M_\Delta^2 &= \frac{1}{2}\mu\frac{v^2}{v_\Delta} \end{aligned} \quad (13)$$

hold, and hence for vevs of the isodoublet and isotriplet, we obtain

$$v^2 = \frac{m_\Phi^2 M_\Delta^2}{\lambda M_\Delta^2 - \mu^2}, \quad (14)$$

$$v_\Delta = \frac{\mu m_\Phi^2}{2\lambda M_\Delta^2 - 2\mu^2} = \frac{\mu}{2}\frac{v^2}{M_\Delta^2}. \quad (15)$$

According to (12), the terms quadratic in φ and δ are

$$V(\varphi,\delta) = \frac{1}{2}m_\Phi^2\varphi^2 + \frac{1}{2}M_\Delta^2\delta^2 - \mu v\varphi\delta. \quad (16)$$

Here and below, the terms suppressed by $(v_\Delta/v)^2$ are omitted.

Denoting the states with definite masses by h and H, we obtain

$$\begin{bmatrix}\varphi\\ \delta\end{bmatrix} = \begin{bmatrix}\cos\alpha & -\sin\alpha\\ \sin\alpha & \cos\alpha\end{bmatrix}\begin{bmatrix}h\\ H\end{bmatrix}, \\ \tan 2\alpha = \frac{2\mu v}{M_\Delta^2 - m_\Phi^2}, \quad (17)$$

$$M_h^2 = \frac{1}{2}\left(m_\Phi^2 + M_\Delta^2 - \sqrt{\left(M_\Delta^2 - m_\Phi^2\right)^2 + 4\mu^2v^2}\right) \approx \\ \approx m_\Phi^2, \quad (18)$$

$$M_H^2 = \frac{1}{2}\left(m_\Phi^2 + M_\Delta^2 + \sqrt{\left(M_\Delta^2 - m_\Phi^2\right)^2 + 4\mu^2v^2}\right) \approx \\ \approx M_\Delta^2. \quad (19)$$

Because $\tan 2\alpha \approx 4v_\Delta/v \ll 1$, the mass eigenstate h consists mostly of φ, and H consists mostly of δ. We suppose that the particle observed by ATLAS and CMS is h, so M_h is about 125 GeV.

The scalar sector of the model, in addition to the massless goldstone bosons that are eaten up by the vector gauge bosons, contains one double-charged field H^{++}, one single-charged field H^+, and three real neutral fields A, H, and h. H^+ is mostly δ^+ with a small Φ^+ admixture, and A is mostly η with a small χ admixture. All these particles except h are heavy; their masses equal M_Δ up to small corrections proportional to v_Δ^2/M_Δ.

2.2. H decays

The second and fourth terms in potential (12) contribute to $H \to 2h$ decays:

$$\frac{\lambda}{2}\left(\Phi^\dagger\Phi\right)^2 \to \frac{\lambda v}{2}\varphi^3, \tag{20}$$

$$\frac{\mu}{\sqrt{2}}\left(\Phi^T i\sigma^2\Delta^\dagger\Phi + \text{H.c.}\right) \to -\frac{\mu}{2}\delta\left(\varphi^2 - \chi^2\right), \tag{21}$$

where in the second line, χ is dominantly a Goldstone state that forms the longitudinal Z polarization.

With the help of (17), we obtain the expression for the effective Lagrangian that describes the $H \to 2h$ decay:

$$\mathcal{L}_{Hhh} = \frac{\mu}{2}\left[1 + \frac{3}{\left(M_H/M_h\right)^2 - 1}\right] Hh^2 = = v_\Delta \frac{M_H^2}{v^2}\left[1 + \frac{3}{\left(M_H/M_h\right)^2 - 1}\right] Hh^2. \tag{22}$$

In the see-saw type-II model, neutrino masses are generated by the Yukawa couplings of the isotriplet Δ to lepton doublets. These couplings generate $H \to \nu\nu$ decays as well. As it was noted in [9], diboson decays dominate for $v_\Delta > 10^{-3}$ GeV. This happens because the diboson decay amplitude is proportional to v_Δ, while Yukawa couplings f_i are inversely proportional to it, $f \sim m_\nu/v_\Delta$. That is why lepton decays are completely negligible for $v_\Delta \gtrsim 1$ GeV.

The amplitudes of $H \to ZZ$ and $H \to W^+W^-$ decays are contained in (7):

$$\mathcal{L}_{HVV} = g^2\left(v_\Delta \cos\alpha - \frac{1}{2}v\sin\alpha\right) W^+W^-H + + \bar{g}^2\left(v_\Delta\cos\alpha - \frac{1}{4}v\sin\alpha\right) Z^2 H \approx \approx -g^2\frac{M_h^2/M_H^2}{1 - M_h^2/M_H^2} v_\Delta W^+W^-H + + \frac{\bar{g}^2}{2}\,\frac{1 - 2M_h^2/M_H^2}{1 - M_h^2/M_H^2} v_\Delta Z^2 H, \tag{23}$$

and we see that the $H \to W^+W^-$ decay is suppressed (see, e. g., [17]).

The $H \to t\bar{t}$ decay occurs through a φ admixture:

$$\mathcal{L}_{Ht\bar{t}} = \sin\alpha\frac{m_t}{v} t\bar{t}H = \frac{2v_\Delta/v}{1 - M_h^2/M_H^2}\,\frac{m_t}{v} t\bar{t}H \tag{24}$$

as well as the H decay into two gluons:

$$\mathcal{L}_{Hgg} = \frac{\alpha_s}{12\pi}\sin\alpha G_{\mu\nu}^2. \tag{25}$$

We note that all the amplitudes of H decays are proportional to the triplet vev v_Δ.

For the decay probabilities, we obtain

$$\Gamma_{H\to hh} = \frac{v_\Delta^2}{v^4}\,\frac{M_H^3}{8\pi}\left[\frac{1 + 2\left(M_h/M_H\right)^2}{1 - \left(M_h/M_H\right)^2}\right]^2 \times \times \sqrt{1 - 4\frac{M_h^2}{M_H^2}}, \tag{26}$$

$$\Gamma_{H\to ZZ} = \frac{v_\Delta^2}{v^4}\,\frac{M_H^3}{8\pi}\left[\frac{1 - 2\left(M_h/M_H\right)^2}{1 - \left(M_h/M_H\right)^2}\right]^2 \times \times \left(1 - 4\frac{M_Z^2}{M_H^2} + 12\frac{M_Z^4}{M_H^4}\right)\sqrt{1 - 4\frac{M_Z^2}{M_H^2}}, \tag{27}$$

$$\Gamma_{H\to WW} = \frac{v_\Delta^2}{v^4}\,\frac{M_H^3}{4\pi}\left[\frac{M_h^2/M_H^2}{1 - \left(M_h/M_H\right)^2}\right]^2 \times \times \left(1 - 4\frac{M_W^2}{M_H^2} + 12\frac{M_W^4}{M_H^4}\right)\sqrt{1 - 4\frac{M_W^2}{M_H^2}}, \tag{28}$$

$$\Gamma_{H\to t\bar{t}} = \frac{v_\Delta^2}{v^4}\,\frac{N_c m_t^2 M_H}{2\pi}\,\frac{1}{\left(1 - M_h^2/M_H^2\right)^2} \times \times \left(1 - 4\frac{m_t^2}{M_H^2}\right)^{3/2}, \tag{29}$$

where $N_c = 3$ is the number of colors. Finally, for the width of decay into two gluon jets, we obtain

$$\Gamma_{H\to gg} = \frac{v_\Delta^2}{v^4}\,\frac{M_H^3}{2\pi}\left(\frac{\alpha_s}{3\pi}\right)^2\left(1 - \frac{M_h^2}{M_H^2}\right)^{-2}, \tag{30}$$

which is always negligible.

In what follows, we suppose that $M_H < 350$ GeV and the decay $H \to t\bar{t}$ is forbidden kinematically. We note that even for $M_H > 350$ GeV, the $H \to 2h$ decay branching ratio is large, but the H production cross section becomes small due to the large H mass.

The H production cross section increases when it becomes lighter, but for $M_H < 250$ GeV, the decay $H \to 2h$ is kinematically forbidden. That is why for numerical estimates we took the value $M_H = 300$ GeV for which $H \to 2h$ and $H \to ZZ$ decays dominate[3)] and $\Gamma_{H\to 2h}/\Gamma_{H\to ZZ} \approx 4$. Hence, a 300-GeV (or slightly lighter) H mostly decays into two 125-GeV Higgs bosons.

3) The decay $H \to ZZ \to \left(l^+l^-\right)\left(l^+l^-\right)$ provides great opportunity for the discovery of a heavy Higgs H.

Table 1. The cross sections of Higgs production via gg fusion. Values for the SM Higgs are taken from Table 4 in [18]. All numbers in this and following tables correspond to 14-TeV LHC energy

M_h, GeV	125	300
$\sigma_{gg\to h}$, pb	$49.97 \pm 10\,\%$	$11.07 \pm 10\,\%$
M_H, GeV	–	300
$\sigma_{gg\to H}$, fb	–	$25 \pm 10\,\%$

A technical remark: the equality $\Gamma_{H\to hh} = \Gamma_{H\to ZZ}$ in the limit $M_H \gg M_h, M_H \gg M_Z$ follows from the equality (up to the sign) of $H \to 2h$ and $H \to 2\chi$ decay amplitudes (see (21)).

2.3. *H* production at the LHC

The dominant mechanism of H production is the gluon fusion, whose cross section equals that of SM Higgs production times $\sin^2\alpha \approx$ $\approx \left[(2v_\Delta/v)\,/\left(1 - M_h^2/M_H^2\right)\right]^2 \approx 2.4 \cdot 10^{-3}$. The relevant numbers are presented in Table 1. All the numbers correspond to the 14-TeV LHC energy.

The subdominant mechanisms of H production are ZZ fusion and associative ZH production. Comparing ZZh and ZZH vertices, we recalculate the cross sections of SM processes of h production into that of H production. In the SM, we have

$$\mathcal{L}_{hZZ} = \frac{1}{4}\bar{g}^2 v Z^2 h. \quad (31)$$

From (23), we obtain

$$\sigma_{ZZ\to H} = \left(\frac{2v_\Delta}{v}\,\frac{1 - 2M_h^2/M_H^2}{1 - M_h^2/M_H^2}\right)^2 \left(\sigma_{ZZ\to h}\right)^{\rm SM} \approx \approx 10^{-3}\left(\sigma_{ZZ\to h}\right)^{\rm SM}, \quad (32)$$

and the same relation holds for the $Z^* \to ZH$ associative production cross section.

We separate the VBF cross section of SM Higgs production into that in W^+W^- fusion (which dominates) and in ZZ fusion (which is the one that matters for H production) with the help of the computer code HAWK [19]. The obtained results are presented in Table 2.

In Table 3, the results for the associative ZH production cross sections are presented.

We see that gluon fusion dominates H production at the LHC. Using model parameters $v_\Delta = 5$ GeV and $M_H = 300$ GeV, we obtain that the branching ratio of $H \to 2h$ decay is $\approx 80\,\%$. Thus, decays of H provide ≈ 20 fb of the double h production cross section in addition to the 40 fb coming from the SM. However, unlike in the SM, where the $2h$ invariant mass is spread along a rather large interval, in the case of H decays the $2h$ invariant mass equals M_H.

Table 2. The cross sections (QCD NLO) of scalar bosons production in vector boson fusion calculated with the help of HAWK (see also Table 10 in [18])

M_h, GeV	125	300
$\sigma_{VV\to h}$, fb	4342(5)	1418(1)
$\sigma_{W^+W^-\to h}$, fb	3272(4)	1053(1)
$\sigma_{ZZ\to h}$, fb	1087(1)	365(1)
M_H, GeV	–	300
$\sigma_{ZZ\to H}$, fb	–	0.365(1)

Table 3. The cross sections of the associative SM Higgs production from Table 14 in [18] and of associative H production recalculated with the help of (32)

M_h, GeV	125	300
$\sigma_{W^*\to Wh}$, fb	$1504 \pm 4\,\%$	$67.6 \pm 4\,\%$
$\sigma_{Z^*\to Zh}$, fb	$883 \pm 5\,\%$	$41.6 \pm 5\,\%$
M_H, GeV	–	300
$\sigma_{Z^*\to ZH}$, fb	–	$0.0416 \pm 5\,\%$

3. *H* PRODUCTION ENHANCEMENT IN THE GEORGI–MACHACEK VARIANT OF THE SEE-SAW TYPE-II MODEL

The amplitudes of H production via both gg fusion and VBF are proportional to the triplet vev v_Δ and because of the upper bound $v_\Delta < 5$ GeV, these amplitudes and the corresponding cross sections are severely suppressed.

The triplet vev v_Δ should be small in order to avoid a noticeable violation of custodial symmetry that guarantees the degeneracy of W and Z bosons in the SM at tree level in the limit $g' = 0, \cos\theta_W = 1$. The vev of the complex isotriplet $\boldsymbol{\Delta}$ with the hypercharge $Y_\Delta = 2$ violates the custodial symmetry (see (8)). The custodial symmetry is preserved when two isotriplets (complex $\boldsymbol{\Delta}$ and real $\boldsymbol{\xi}$ with $Y_\xi = 0$) are added to the SM and when vevs of their neutral components are equal [12]. Thus, in the GM variant of the see-saw type-II model, v_Δ is not bounded by (10) and can be considerably larger. Instead of (8), in the GM model we have

$$M_W^2 = \frac{g^2}{4}\left(v^2 + 4v_\Delta^2\right), \quad M_Z^2 = \frac{\bar{g}^2}{4}\left(v^2 + 4v_\Delta^2\right), \tag{33}$$

and instead of (11),

$$v^2 + 4v_\Delta^2 = (246\ \text{GeV})^2. \tag{34}$$

We note that our v_Δ is $\sqrt{2}$ times larger than what is usually used in the papers devoted to the GM model; our v is also usually denoted by v_Φ, while the value 246 GeV is denoted by v.

The scalar particles are conveniently classified in the GM model by their transformation properties under the custodial $SU(2)$. Two singlets which mix to form mass eigenstates h and H are

$$H_1^0 = \varphi, \quad H_2^0 = \sqrt{\frac{2}{3}}\delta + \sqrt{\frac{1}{3}}\xi^0 \tag{35}$$

(see, e. g., [20]). Due to the considerable admixture of ξ^0 in H_2^0, the HW^+W^- coupling constant is not suppressed and three modes of H decays are essential: $H \to hh$, $H \to W^+W^-$, and $H \to ZZ$.

The recently discovered Higgs boson should be identified with h. The deviations of h couplings to vector bosons and fermions from their values in the SM lead to the upper bound on v_Δ. These deviations in the limit of heavy scalar triplets were recently studied in [20] (see also [21]). From Eqs. (59) and (61) in [20], we obtain the following estimates for the ratios of the hVV (here $V = W,\ Z$) and $h\bar{f}f$ coupling constants to that in the SM:

$$k_V \approx 1 + 3\left(\frac{v_\Delta}{v}\right)^2, \quad k_f \approx 1 - \left(\frac{v_\Delta}{v}\right)^2. \tag{36}$$

Because the Higgs boson h is produced at the LHC mainly in gluon fusion through a t-quark triangle, we obtain the ratio of the cross sections to that in the SM as

$$\mu_{\tau\bar{\tau}} \approx 1 - \left(2\frac{v_\Delta}{v}\right)^2, \quad \mu_{VV} \approx 1 + \left(2\frac{v_\Delta}{v}\right)^2. \tag{37}$$

Since the $h \to b\bar{b}$ decay is studied in associative production, $V^* \to Vh \to Vb\bar{b}$, we have

$$\mu_{b\bar{b}} \approx 1 + \left(\frac{2v_\Delta}{v}\right)^2. \tag{38}$$

In (37) and (38) we used that the total width of h is practically the same as in the SM.

Finally, in the case of the $h \to \gamma\gamma$ decay, the SM factor $16/9 - 7$ in the amplitude is modified as

$$\frac{16}{9} - 7 \to \left[1 - \left(\frac{v_\Delta}{v}\right)^2\right]\left[\frac{16}{9}\left(1 - \left(\frac{v_\Delta}{v}\right)^2\right) - 7\left(1 + 3\left(\frac{v_\Delta}{v}\right)^2\right)\right] = \frac{16}{9}\left(1 - 2\left(\frac{v_\Delta}{v}\right)^2\right) - 7\left(1 + 2\left(\frac{v_\Delta}{v}\right)^2\right), \tag{39}$$

where the first factor in the first line takes the damping of h production in gluon fusion into account[4)].

We suppose that v_Δ is ten times larger than the number used in Sec. 2, $v_\Delta^{GM} = 50$ GeV. Then from (34) we obtain $v^{GM} \approx 225$ GeV and $\mu_{\tau\bar{\tau}} \approx 0.8$, while $\mu_{WW} = \mu_{ZZ} = \mu_{bb} \approx 1.2$. From (39), $\mu_{\gamma\gamma} \approx 1.4$. With the up-to-date level of experimental accuracy, one cannot exclude these deviations of the μ_i from their SM values $(\mu_i)^{SM} \equiv 1$.

One order of magnitude increase in v_Δ leads to two orders of magnitude increase in the H production cross section. Hence, a 300 GeV heavy Higgs boson H can be produced at the 14-TeV LHC energy with the 2 pb cross section, which should be large enough for it to be discovered prior to the HL-LHC. The search strategy should be the same as for the SM Higgs boson: the $gg \to H \to ZZ$ decay is a golden discovery mode, whose cross section can be as large as $(2\ \text{pb}) \times \text{Br}\,(H \to ZZ)^{\text{GM}}$, where $\text{Br}\,(H \to ZZ)^{\text{GM}}$ depends on the model parameters (see [20]).

4. CONCLUSIONS

The case of extra isotriplet(s) provides a rich Higgs-sector phenomenology with charged and neutral scalar particles additional to the SM Higgs boson. With the growth of the triplet vev, the production cross section of the new scalar grows and the dominant decays of new particles become decays to gauge and lighter scalar bosons. The charged scalars (Φ^{++}, Φ^{+}) are produced through electroweak interactions. The bounds on the model parameters from the nondiscovery of Φ^{++} and

4) We take only t-quark and W-boson loops into account, omitting the loops with charged Higgses.

Φ^+ with the 8-TeV LHC data and the prospects of their discovery at the 14-TeV LHC energy are discussed, in particular, in [22]. In this paper, we have discussed the neutral heavy Higgs production at the LHC in which the gluon fusion dominates. The $H \to 2h$ decay contributes significantly to the double Higgs production and even may dominate in the GM variant of the see-saw type-II model. The best discovery mode for H is the "golden mode" $pp \to HX \to ZZX$, and its cross section can be only a few times smaller than for the heavy SM Higgs.

After this paper had been written, paper [23] appeared in arXiv in which the enhancement of double Higgs production due to heavy Higgs decay is considered in the framework of the MSSM model with two isodoublets. The $H \to 2h$ resonant decay in the MSSM at small $\tan\beta$ was previously analyzed in [10].

This paper is our present to Valery Anatolievich Rubakov on his anniversary. Many students (and not only students) in the world are studying Physics by reading his excellent books and papers and listening to his brilliant lectures.

We are grateful to A. Denner for the clarifications concerning HAWK code and to J. Baglio, C. Englert, G. Senjanović, and J. Zurita for providing us with relevant references. The authors are partially supported by the RFBR (grants Nos. 12-02-00193, 14-02-00995), and the grant NSh-3830.2014.2. One of the authors (S. G.) was also supported by Dynasty Foundation and by the Russian Federation Government under grant No. 11.G34.31.0047.

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