Search for the Higgs boson H_2^0 at LHC in 3-3-1 model

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We present an analysis of production and signature of neutral Higgs boson (H_2^0) on the version of the 3-3-1 model containing heavy leptons at the Large Hadron Collider. We studied the possibility to identify it using the respective branching ratios. Cross section are given for the collider energy, $\sqrt{s} = 14$ TeV. Event rates and significances are discussed for two possible values of integrated luminosity, 300 fb^{-1} and 3000 fb^{-1}.

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I. INTRODUCTION

The way to the understanding of the symmetry breaking mainly go through the scalars, although there are many other models that do not contain elementary scalar fields, such as the Nambu–Jona–Lasinio mechanism, technicolor theories, and the strongly interacting gauge systems [1]. These scalars protect the renormalizability of the theory by moderating the cross section growth. But so far, despite many experimental and theoretical efforts in order to understand the scalar sector, the Higgs mechanism remains still unintelligible. Nowadays, the major goal of the experimentalists in particle physics at the LHC is to unravel the nature of electroweak symmetry breaking. The standard model (SM) is the prototype of a gauge theory with spontaneous symmetry breaking. This had great success in explaining most of the experimental data. However, recent results from neutrino oscillation experiments make clear that the SM is not complete; then, the neutrino oscillation implies that at least two neutrino flavors are massive. Moreover, there are other crucial problems in particle physics that do not get a response in the SM. For instance, it offers no solution to the dark matter problem, dark energy, or the asymmetry of matter-antimatter in the Universe. Therefore, there is a consensus among particle physicists that the SM must be extended.

In the SM, there appears only one elementary scalar, which arises through the breaking of electroweak symmetry, and this is the Higgs boson. The Higgs boson is an important prediction of several quantum field theories and is so crucial to our understanding of the Universe. So on July 4, 2012, a previously unknown particle with a mass around 126 GeV was announced as being detected, which physicists suspected at the time to be the Higgs boson [2–4]. By March 2013, the particle had been proven to be the Higgs boson because it behaves, interacts, and decays in the ways predicted by the SM and was also tentatively confirmed to have positive parity and zero spin, two fundamental criteria of a Higgs boson, making it also the first known scalar particle to be discovered in nature.

Different types of Higgs bosons, if they exist, may lead us into new realms of physics beyond the SM. Since the SM leaves many questions open, there are several extensions. For example, if the grand unified theory (GUT) contains the SM at high energies, then the Higgs bosons associated with GUT symmetry breaking must have masses of order $M_X \sim \mathcal{O}(10^{15})$ GeV. Supersymmetry [5] provides a solution to the hierarchy problem through the cancellation of the quadratic divergences via fermionic and bosonic loops contributions [6]. Moreover, the minimal supersymmetric extension of the SM can be derived as an effective theory of supersymmetric GUT [7].

Among these extensions of the SM, there are also other classes of models based on $\mathrm{SU}(3)_C \otimes \mathrm{SU}(3)_L \otimes \mathrm{U}(1)_N$ gauge symmetry (3-3-1 model) [8–10], in which the anomaly cancellation mechanisms occur when the three basic fermion families are considered and not family by family as in the SM. This mechanism is peculiar because it requires that the number of families is an integer multiple of the number of colors. This feature combined together with the asymptotic freedom, which is a property of quantum chromodynamics, requires that the number of families is 3. Moreover, according to these models, the Weinberg angle is restricted to the value $s_W^2 = \sin^2\theta_W < 1/4$ in the version of heavy leptons [8]. Thus, when it evolves to

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higher values, it shows that the model loses its perturbative character when it reaches a mass scale of about 4 TeV [11]. Hence, the 3-3-1 model is one of the most interesting extensions of the SM and is phenomenologically well motivated to be probed at the LHC and other accelerators.

In this work, we study the production and signatures of an extended neutral Higgs boson H_2^0, predicted by the 3-3-1 model, which incorporates the charged heavy leptons [8,12]. We can show that the neutral Higgs boson signatures can be significant at the LHC. The signal of the new particle can be obtained by studying the different decay modes and if we consider a luminosity of 10 times higher than the original LHC design. With respect to both mechanisms, that is, the Drell–Yan and gluon-gluon fusion, we consider the Z', H_1^0, and H_2^0 as propagators. Therefore, in Sec. II, we present the relevant features of the model. In Sec. III, we compute the total cross sections of the process, and in Sec. IV, we summarize our results and conclusions.

II. RELEVANT FEATURES OF THE MODEL

We are working here with the version of the 3-3-1 model that contains heavy leptons [8]. The model is based on the semisimple symmetry group $\mathrm{SU}(3)_C \otimes \mathrm{SU}(2)_L \otimes \mathrm{U}(1)_N$. The electric charge operator is given by

$$\frac{Q}{e} = (T_3 - \sqrt{3}T_8) + N, \tag{1}$$

where T_3 and T_8 are the generators of SU(3) and e is the elementary electric charge. So, we can build three triplets of quarks of $\mathrm{SU}(3)_L$ such that

$$\begin{aligned} Q_{1L} &= \begin{pmatrix} u_1' \\ d_1' \\ J_1 \end{pmatrix}_L \sim \left(\mathbf{3}, \frac{2}{3}\right), \\ Q_{\alpha L} &= \begin{pmatrix} J_\alpha' \\ u_\alpha' \\ d_\alpha' \end{pmatrix}_L \sim \left(\mathbf{3}^*, -\frac{1}{3}\right), \end{aligned} \tag{2}$$

where the new quark J_1 carries 5/3 units of electric charge while J_α ($\alpha = 2, 3$) carry $-4/3$ each. We must also introduce the right-handed fermionic fields $U_R \sim (\mathbf{1}, 2/3)$, $D_R \sim (\mathbf{1}, -1/3)$, $J_{1R} \sim (\mathbf{1}, 5/3)$, and $J'_{\alpha R} \sim (\mathbf{1}, -4/3)$. We have defined

$$U = \begin{pmatrix} u' & c' & t' \end{pmatrix}$$

and

$$D = \begin{pmatrix} d' & s' & b' \end{pmatrix}.$$

The spontaneous symmetry breaking is accomplished via three SU(3) scalar triplets, which are

$$\begin{aligned} \eta &= \begin{pmatrix} \eta^0 \\ \eta_1^- \\ \eta_2^+ \end{pmatrix} \sim (\mathbf{3}, 0), \qquad \rho = \begin{pmatrix} \rho^+ \\ \rho^0 \\ \rho^{++} \end{pmatrix} \sim (\mathbf{3}, 1), \\ \chi &= \begin{pmatrix} \chi^- \\ \chi^{--} \\ \chi^0 \end{pmatrix} \sim (\mathbf{3}, -1). \end{aligned} \tag{3}$$

For the sake of simplicity, we will assume here that the model respects the $B + L$ symmetry, where B is the baryon number and L is the lepton number. Then, the more general renormalizable Higgs potential is given by

$$\begin{aligned} V(\eta,\rho,\chi) = {} & \mu_1^2\eta^\dagger\eta + \mu_2^2\rho^\dagger\rho + \mu_3^2\chi^\dagger\chi + \lambda_1(\eta^\dagger\eta)^2 \\ & + \lambda_2(\rho^\dagger\rho)^2 + \lambda_3(\chi^\dagger\chi)^2 + \eta^\dagger\eta[\lambda_4(\rho^\dagger\rho) \\ & + \lambda_5(\chi^\dagger\chi)] + \lambda_6(\rho^\dagger\rho)(\chi^\dagger\chi) + \lambda_7(\rho^\dagger\eta)(\eta^\dagger\rho) \\ & + \lambda_8(\chi^\dagger\eta)(\eta^\dagger\chi) + \lambda_9(\rho^\dagger\chi)(\chi^\dagger\rho) \\ & + \frac{1}{2}(f\varepsilon^{ijk}\eta_i\rho_j\chi_k + \mathrm{c.H.}), \end{aligned} \tag{4}$$

where μ_i ($i = 1, 2, 3$) and f are constants with mass dimension and λ_j ($j = 1, \ldots, 9$) are dimensionless constants [12]. The potential (4) is bounded from below when the neutral Higgs fields develop the vacuum expectation values (VEVs) $\langle\eta^0\rangle = v_\eta$, $\langle\rho^0\rangle = v_\rho$ and $\langle\chi^0\rangle = v_\chi$, with $v_\eta^2 + v_\rho^2 = v_W^2 = 246^2\ \mathrm{GeV}^2$. The scalar χ^0 is supposedly heavy, and it is responsible for the spontaneous symmetry breaking of $\mathrm{SU}(3)_L \otimes \mathrm{U}(1)_N$ to $\mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y$ of the standard model. Meanwhile, η^0 and ρ^0 are lightweight and are responsible for the breaking of $\mathrm{SU}(2)_L \otimes \mathrm{U}(1)_Y$ to $\mathrm{U}(1)_Q$ of the electromagnetism. Therefore, it is reasonable to expect

$$v_\chi \gg v_\eta, v_\rho. \tag{5}$$

The potential (4) provides the masses of the neutral Higgs as

$$m_{H_1^0}^2 \approx 4\frac{\lambda_2 v_\rho^4 - \lambda_1 v_\eta^4}{v_\eta^2 - v_\rho^2}, \qquad m_{H_2^0}^2 \approx \frac{v_W^2 v_\chi^2}{2v_\eta v_\rho}, \tag{6a}$$

$$m_{H_3^0}^2 \approx -\lambda_3 v_\chi, \qquad m_h^2 = -\frac{f v_\chi}{v_\eta v_\rho}\left[v_W^2 + \left(\frac{v_\eta v_\rho}{v_\chi}\right)^2\right], \tag{6b}$$

with the corresponding eigenstates

$$\begin{pmatrix} \xi_\eta \\ \xi_\rho \end{pmatrix} \approx \begin{pmatrix} c_w & s_w \\ s_w & c_w \end{pmatrix}\begin{pmatrix} H_1^0 \\ H_2^0 \end{pmatrix}, \qquad \xi_\chi \approx H_3^0, \qquad \zeta_\chi \approx ih, \tag{7}$$

where the mixing parameters are $c_w = \cos w = v_\eta/\sqrt{v_\eta^2 + v_\rho^2}$ and $s_w = \sin w$ [12]. In Eqs. (6) and (7), we used the approximation (5), and, so as not to introduce the new mass scale in the model, we assume $f \approx -v_\chi$. We can then note that H_3^0 is a typical 3-3-1 Higgs boson. The scalar H_1^0 is one that can be identified with the SM Higgs since its mass and eigenstate do not depend on v_χ.

Now, we can write the Yukawa interactions for the ordinary quarks, i.e.,

$$\mathcal{L}_q^Y = \sum_\alpha \Big[\bar{Q}_{1L}(G_{1\alpha} U'_{\alpha R}\eta + \tilde{G}_{1\alpha} D'_{\alpha R}\rho) + \sum_i \bar{Q}_{iL}(F_{i\alpha} U'_{\alpha R}\rho^* + \tilde{F}_{i\alpha} D'_{\alpha R}\eta^*) \Big], \tag{8}$$

where G_{ab} and G'_{ab} (a and b are generation indexes) are coupling constants.

The interaction eigenstates (2) and their right-handed counterparts can rotate about their respective physical eigenstates as

$$U'_{aL(R)} = \mathcal{U}_{ab}^{L(R)} U_{bL(R)}, \tag{9a}$$

$$D'_{aL(R)} = \mathcal{D}_{ab}^{L(R)} U_{bL(R)}, \qquad J'_{aL(R)} = \mathcal{J}_{ab}^{L(R)} J_{bL(R)}. \tag{9b}$$

Since the cross sections involving the sum over the flavors and rotation matrices are unitary, then the mixing parameters have no major effect on the calculations. In terms of physical fields, the Yukawa Lagrangian for the neutral Higgs can be written as

$$\begin{aligned} -\mathcal{L}_Q = \frac{1}{2}\Big\{ & \bar{U}(1+\gamma_5)\Big[1 + \Big[\frac{s_w}{v_\rho} + \Big(\frac{c_w}{v_\eta} + \frac{s_w}{v_\rho}\Big)\mathcal{V}^U\Big]H_1^0 \\ & + \Big[\frac{-c_w}{v_\rho} + \Big(\frac{s_w}{v_\eta} - \frac{c_w}{v_\rho}\Big)\mathcal{V}^U\Big]H_2^0\Big]M^U U \\ & + \bar{D}(1+\gamma_5)\Big[1 + \Big[\frac{c_w}{v_\eta} + \Big(\frac{s_w}{v_\rho} - \frac{c_w}{v_\eta}\Big)\mathcal{V}^D\Big]H_1^0 \\ & + \Big[\frac{s_w}{v_\eta} - \Big(\frac{c_w}{v_\rho} + \frac{s_w}{v_\eta}\Big)\mathcal{V}^D\Big]H_2^0\Big]M^D D\Big\} + \text{H.c.}, \end{aligned} \tag{10}$$

where $V_L^U V_L^D = V_{\text{CKM}}$ is the Cabibbo–Kobayashi–Maskawa matrix, $\mathcal{V}^U$ and $\mathcal{V}^D$ are arbitrary mixing matrices, and

$$M^U = \text{diag}(m_u \quad m_c \quad m_t)$$

and

$$M^D = \text{diag}(m_d \quad m_s \quad m_b)$$

are matrices that carry the masses of the quarks.

In the gauge sector, beyond the standard particles γ, Z, and $W^\pm$, the model predicts one neutral (Z'), two single-charged ($V^\pm$), and two doubly-charged ($U^{\pm\pm}$) gauge bosons. The gauge interactions with Higgs bosons are given by

$$\mathcal{L}_{GH} = \sum_\varphi (\mathcal{D}_\mu \varphi)^\dagger (\mathcal{D}_\mu \varphi), \tag{11}$$

where the covariant derivatives are

$$\mathcal{D}_\mu \varphi_i = \partial_\mu \varphi_i - ig\Big(W_\mu . \frac{T}{2}\Big)_i^j \varphi_j - ig' N_\varphi \varphi_i B_\mu, \tag{12}$$

where $\varphi = \eta, \rho, \chi$ ($N_\eta = 0$, $N_\rho = 1$, $N_\chi = -1$) are the Higgs triplets, W_μ and B_μ are the SU(2) and U(1) field tensors, and g and g' are the U(1) and SU(2) coupling constants, respectively. Diagonalization of the Lagrangean (11), after symmetry breaking, gives masses for the neutral weak gauge bosons, i.e.,

$$m_Z \approx \frac{|e|}{2 s_W c_W} v_W, \qquad m_{Z'}^2 \approx \frac{1}{3(1-4s_W^2)}\Big(\frac{|e| c_W v_\chi}{s_W}\Big)^2, \tag{13}$$

where $s_W = \sin\theta_W$, with θ_W being the Weinberg angle, and $c_W^2 = 1 - s_W^2$. Then, the eigenstates are

$$W_\mu^3 \approx s_W A_\mu - c_W Z_\mu \tag{14a}$$

$$W_\mu^8 \approx -\sqrt{3} s_W\Big(A_\mu - \frac{s_W}{c_W} Z_\mu\Big) + \frac{\sqrt{1-4s_W^2}}{c_W} Z'_\mu \tag{14b}$$

$$B_\mu \approx \frac{s_W}{\sqrt{1-4s_W^2}} A_\mu + \frac{s_W}{c_W}\Big(Z_\mu + \sqrt{3} Z'_\mu\Big). \tag{14c}$$

In Eqs. (13) and (14), we have used the approximation (5). Finally, the weak neutral current in the sector of u and d quarks reads

$$-\mathcal{L}_Z = \frac{|e|}{2 s_W c_W} \bar{q}\gamma^\mu [v(q) + a(q)\gamma_5] q Z_\mu \tag{15a}$$

$$-\mathcal{L}_{Z'} = \frac{|e|}{2 s_W c_W} \bar{q}\gamma^\mu [v'(q) + a'(q)\gamma_5] q Z'_\mu, \tag{15b}$$

for which the coefficients are

$$v(u) = 1 - \frac{s_W^2}{8}, \qquad a(u) = -a(d) = -1,$$

$$v(d) = -1 + \frac{4}{3} s_W^2, \tag{16a}$$

$$v'(u) = \sqrt{1 + 4s_W^2}, \qquad a'(u) = \sqrt{\frac{1-4s_W^2}{3}},$$

$$v'(d) = \frac{2s_W^2 - 1}{\sqrt{3}}, \tag{16b}$$

$$a'(d) = -v'(d). \tag{16c}$$

In this work, we study the production of the neutral Higgs boson H_2^0 at pp colliders. With respect to both mechanisms, that is, the Drell–Yan and gluon-gluon fusion, we consider the Z', H_1^0, and H_2^0 as propagators.

III. CROSS SECTION PRODUCTION

The mechanisms for the production of a neutral Higgs particle H_2^0 in pp collisions occur in association with the bosons Z', H_1^0, and H_2^0; see Figs. 1 and 2. Unlike of the

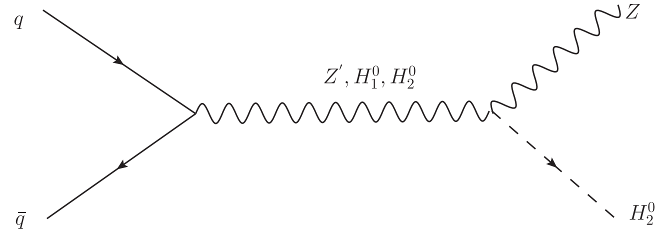

FIG. 1. Feynman diagrams for the production of a neutral Higgs via the Drell–Yan process.

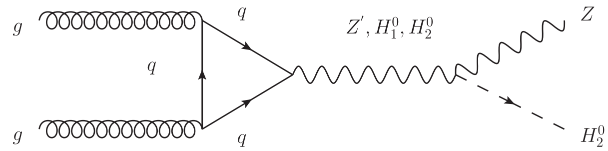

FIG. 2. Feynman diagrams for the production of a neutral Higgs via gluon-gluon fusion.

SM, where the gluon-gluon fusion dominates over Drell–Yan when the Higgs boson is heavier than 100 GeV [13], in 3-3-1 Model does not occur. Here, the mechanism of Drell–Yan dominates over gluon-gluon fusion at leading order for H_2^0 production at $\sqrt{s} = 14$ TeV. The process $pp \to H_2^0 Z$ ($i = 1, 2$) takes place in the s channel. The term involving the boson Z is absent because there is no coupling between the Z and $H_2^0 Z$; moreover, the interference term between the Z' and H_2^0 should be absent because it gives an imaginary value. So, using the interaction Lagrangian [8,14], we obtain the differential cross section in the first place for Drell–Yan for H_2^0,

$$\frac{d\hat{\sigma}_{H_2^0}}{d\cos\theta} = \frac{\beta_{H_2^0} g^2}{192\pi c_W^2 s}\Bigg\{\frac{g^4\Lambda^2_{ZZ'H_2^0}}{4(s-m_{Z'}^2+im_{Z'}\Gamma_{Z'})^2}\Bigg(\Bigg(m_Z^2+\frac{tu}{m_Z^2}-t-u+s\Bigg)(g^2_{V'^q}+g^2_{A'^q})\Bigg)$$
$$+\Bigg(\frac{v_\eta^2 v_\rho^2 m_q^2}{2v_W^6}|\chi^{(1)}(\hat{s})|^2+\frac{(m_u\frac{v_\eta}{v_\rho}-m_d\frac{v_\rho}{v_\eta})^2(v_\rho^2-v_\eta^2)^2}{32v_W^6}|\chi^{(2)}(\hat{s})|^2+\frac{m_q(m_u\frac{v_\eta}{v_\rho}-m_d\frac{v_\rho}{v_\eta})v_\eta v_\rho(v_\rho^2-v_\eta^2)}{4v_W^6}|\chi^{(1)}(\hat{s})||\chi^{(2)}(\hat{s})|\Bigg)$$
$$\times\Bigg(\frac{\hat{s}}{m_Z^2}(\hat{s}^2-2m_Z^2\hat{s}+m_Z^4)-\frac{m_q^2}{m_Z^2}(2\hat{s}^2-4\hat{s}m_Z^2+2m_Z^4)+-\frac{m^2_{H_2^0}}{m_Z^2}(2\hat{s}^2+2\hat{s}m_q^2-4m_q^2\hat{s}-4m_q^2m_Z^2)+\frac{m^4_{H_2^0}\hat{s}}{m_Z^2}-\frac{2m^4_{H_2^0}m_q^2}{m_Z^2}\Bigg)\Bigg\}; \tag{17}$$

here, g is the coupling constant of the weak interaction, $\beta_{H_2^0}$ is the Higgs velocity in the c.m. of the subprocess, which is equal to

$$\beta_{H_2^0} = \frac{\Big[\Big(1-\frac{(m_Z+m_{H_2^0})^2}{\hat{s}}\Big)\Big(1-\frac{(m_Z-m_{H_2^0})^2}{\hat{s}}\Big)\Big]^{1/2}}{1-\frac{m_Z^2-m^2_{H_2^0}}{\hat{s}}},$$

and t and u are

$$t = m_q^2+m_Z^2-\frac{s}{2}\Bigg\{\Bigg(1+\frac{m_Z^2-m_H^2}{s}\Bigg) - \cos\theta\Bigg[\Bigg(1-\frac{(m_Z+m_H)^2}{s}\Bigg)\Bigg(1-\frac{(m_Z-m_H)^2}{s}\Bigg)\Bigg]^{1/2}\Bigg\},$$
$$u = m_q^2+m_H^2-\frac{s}{2}\Bigg\{\Bigg(1-\frac{m_Z^2-m_H^2}{s}\Bigg) + \cos\theta\Bigg[\Bigg(1-\frac{(m_Z+m_H)^2}{s}\Bigg)\Bigg(1-\frac{(m_Z-m_H)^2}{s}\Bigg)\Bigg]^{1/2}\Bigg\},$$

where θ is the angle between the Higgs and the incident quark in the c.m. frame. We have also defined

$$\chi^i(\hat{s}) = \frac{1}{\hat{s}-m^2_{H_i^0}+im_{H_i^0}\Gamma_{H_i^0}},$$

with $\Gamma_{H_i^0}$ being the Higgs boson total width; $i = 1, 2$, $\Gamma_{Z'}$ [14,15] are the total width of the Z' boson, m_q, where $q = u$, d are the masses of the quark; $g^q_{V',A'}$ are the 3-3-1 quark coupling constants; $\sqrt{\hat{s}}$ is the center of mass energy of the $q\bar{q}$ system; $g = \sqrt{4\pi\alpha}/\sin\theta_W$; and α is the fine structure constant, which we take equal to $\alpha = 1/128$. For the Z' boson, we take $M_{Z'} = (0.5–3)$ TeV, since $M_{Z'}$ is proportional to the VEV v_χ [9,10]. For the standard model parameters, we assume Particle Data Group values, i.e., $M_Z = 91.19$ GeV, $\sin^2\theta_W = 0.2315$, and $M_W = 80.33$ GeV [16], and t and u are the kinematic invariants. We have also defined the $\Lambda_{ZZ'H_2^0}$ as the coupling constants of the Z' boson to the Z boson and Higgs H_2^0; the $\Lambda_{H_i^0 H_2^0 Z}$ are the couplings constants of the H_1^0 boson to H_2^0 and the Z boson and of the H_2^0 boson to H_2^0 and the Z boson. These coupling constants should be multiplied by $p^\mu - q^\mu$ to get a $\Lambda^\mu_{H_i^0 H_2^0 Z} = \Lambda_{H_i^0 H_2^0 Z}(p^\mu - q^\mu)$ with p and q being the momentum 4-vectors of the H_2 and Z boson, and the $\Lambda_{q\bar{q}H_i^0}$ are the coupling constants of the $H_1^0(H_2^0)$ to $q\bar{q}$; the $v(q)$, $a(q)$, $v'(q)$ and $a'(q)$ are given in Ref. [15]. We remark still that, in 3-3-1 model, the states H_1^0 and H_2^0 are mixed:

$$(\Lambda_{q\bar{q}Z'})_\mu \approx i\frac{|e|}{2s_W c_W}\gamma_\mu[v'(q)+a'(q)\gamma_5], \tag{18a}$$

$$\Lambda_{q\bar{q}H_1^0} \approx -i\frac{m_q}{2v_W}(1+\gamma_5), \tag{18b}$$

$$\Lambda_{q\bar{q}H_2^0} \approx \frac{i}{2v_W}\left(-m_u\frac{v_\eta}{2v_\rho}+m_d\frac{v_\rho}{v_\eta}\right)(1+\gamma_5) \tag{18c}$$

$$(\Lambda_{ZZ'H_2^0})_{\mu\nu} \approx \frac{g^2}{\sqrt{3}(1-4s_W^2)}\frac{v_\eta v_\rho}{v_W}g_{\mu\nu}, \tag{18d}$$

$$(\Lambda_{H_2^0H_2^0Z})_\mu \approx -\frac{g}{2}\frac{m_Z}{m_W}\frac{(v_\rho^2-v_\eta^2)}{v_W^2}(p-q)_\mu, \tag{18e}$$

$$(\Lambda_{H_1^0H_2^0Z})_\mu \approx -2g\frac{m_Z}{m_W}\frac{v_\rho v_\eta}{v_W^2}(p-q)_\mu. \tag{18f}$$

The total cross section for the process $pp \to qq \to ZH_2^0$ is related to the subprocess $qq \to ZH_2^0$ total cross section $\hat{\sigma}$ through

$$\sigma = \int_{\tau_{\min}}^{1}\int_{\ln\sqrt{\tau_{\min}}}^{-\ln\sqrt{\tau_{\min}}} d\tau dy q(\sqrt{\tau}e^{y},Q^2)q(\sqrt{\tau}e^{-y},Q^2)\hat{\sigma}(\tau,s), \tag{19}$$

where $\tau_{\min} = (m_Z + m_{H_2^0})^2/s(\tau = \hat{s}/s)$ and $q(x, Q^2)$ is the quark structure function.

Another form to produce a neutral Higgs is via the gluon-gluon fusion, namely, through the reaction of the type $pp \to gg \to ZH_2^0$. Since the final state is neutral, the s channel involves the exchange of the boson Z', H_1^0, and H_2^0. The exchange of a photon is not allowed by C conservation (Furry's theorem), which also indicates that only the axial-vector coupling of the boson Z' contributes to this process. It is important to emphasis that, for production of H_2^0, we take the interference between the Z', which are antisymmetric in the gluon polarizations, and the H_2^0 (we only consider the antisymmetric term of H_2^0) because the other part is symmetric and therefore vanishes; then, we write explicitly the Z', H_1^0, and H_2^0 contributions to the elementary cross section for the production of H_2^0,

$$\left(\frac{d\hat{\sigma}}{d\cos\theta}\right)^{Z'}_{pp\to ZH_2^0} = \frac{g^6\alpha_s^2(\Lambda_{Z(Z')ZH_i^0})^2\Delta}{8192\pi^3\hat{s}c_W^2M_{Z(Z')}^4}\beta_{H_i^0} \times \left|\sum_{q=u,d}T_3^q(q')(1+2\delta_q I_q)\right|^2, \tag{20}$$

$$\left(\frac{d\hat{\sigma}}{d\cos\theta}\right)^{H_1^0-H_2^0}_{pp\to ZH_2^0} = \frac{g^2\alpha_s^2(v_\rho^2-v_\eta^2)v_\rho v_\eta\Omega\beta_{H_i^0}}{8192\pi^3\hat{s}c_W^2v_W^6}\mathrm{Re}\chi^{(1)}(\hat{s})\chi^{(2)}(\hat{s})\sum_{q=u,d}m_q^3\left(m_u\frac{v_\eta}{v_\rho}-m_d\frac{v_\rho}{v_\eta}\right)I_q\sum_{q=u,d}I_q^* + \frac{g^2\alpha_s^2(v_\rho^2-v_\eta^2)v_\rho v_\eta\hat{s}\Omega\beta_{H_i^0}}{16384\pi^3c_W^2v_W^6}\mathrm{Re}\chi^{(1)}(\hat{s})\chi^{(2)}(\hat{s})\sum_{q=u,d}\frac{(m_u\frac{v_\eta}{v_\rho}-m_d\frac{v_\rho}{v_\eta})}{m_q}I_q\sum_{q=u,d}I_q^*, \tag{21}$$

$$\left(\frac{d\hat{\sigma}}{d\cos\theta}\right)^{H_1^0}_{pp\to ZH_2^0} = \frac{g^2\alpha_s^2v_\eta^2v_\rho^2\hat{s}\Omega\beta_{H_i^0}}{8192\pi^3v_W^6c_W^2}|\chi^{(1)}(\hat{s})|^2\left|\sum_{q=u,d}2\delta_q+\delta_q(4\delta_q-1)I_q\right|^2 + \frac{g^2\alpha_s^2v_\eta^2v_\rho^2\Omega\beta_{H_i^0}}{4096\pi^3\hat{s}c_W^2v_W^6}|\chi^{(1)}(\hat{s})|^2\left|\sum_{q=u,d}m_q^2I_q\right|^2, \tag{22}$$

$$\left(\frac{d\hat{\sigma}}{d\cos\theta}\right)^{H_2^0}_{pp\to ZH_2^0} = \frac{g^2\alpha_s^2(v_\rho^2-v_\eta^2)^2\hat{s}\Omega\beta_{H_i^0}}{131072\pi^3v_W^6c_W^2}|\chi^{(2)}(\hat{s})|^2\left|\sum_{q=u,d}\frac{(m_u\frac{v_\eta}{v_\rho}-m_d\frac{v_\rho}{v_\eta})}{m_q}[2\delta_q+\delta_q(4\delta_q-1)I_q]\right|^2 + \frac{g^2\alpha_s^2(v_\rho^2-v_\eta^2)^2\Omega\beta_{H_i^0}}{65536\pi^3\hat{s}v_W^6c_W^2}|\chi^{(2)}(\hat{s})|^2\left|\sum_{q=u,d}m_q\left(m_u\frac{v_\eta}{v_\rho}-m_d\frac{v_\rho}{v_\eta}\right)I_q\right|^2, \tag{23}$$

$$\left(\frac{d\hat{\sigma}}{d\cos\theta}\right)^{Z'-H_1^0}_{pp\to ZH_2^0} = -\frac{g^4\alpha_s^2\Lambda_{ZZ'H_2^0}v_\eta v_\rho\Pi\beta_{H_i^0}}{1024\pi^3\hat{s}^2c_W^2v_W^3}\mathrm{Re}\left[\frac{\chi^{(1)}(\hat{s})}{(s-m_{Z'}^2+im_{Z'}\Gamma_{Z'})}\sum_{q=u,d}m_q^2T_3^q(1+2\delta_qI_q)\sum_{q=u,d}I_q^*\right], \tag{24}$$

TABLE I. Values for the particle masses used in this work. All the values in this table are given in GeV. Here, $m_{H^{\pm\pm}} = 500$ GeV and $m_T = 2v_\chi$.

f	v_χ, m_{J_1}	m_E	m_M	$m_{H_3^0}$	m_{h^0}	$m_{H_1^0}$	$m_{H_2^0}$	$m_{H_2^\pm}$	m_V	m_U	$m_{Z'}$	$m_{J_{2,3}}$
−1008.3	1000	148.9	875	2000	1454.6	126	1017.2	183	467.5	464	1707.6	1410
−1499.7	1500	223.3	1312.5	474.34	2164.32	125.12	1525.8	387.23	694.12	691.76	2561.3	2115
−1993.0	2000	297.8	1750	632.45	2877.07	125.12	2034.37	519.39	922.12	920.35	3415.12	2820

$$\left(\frac{d\hat{\sigma}}{d\cos\theta}\right)^{Z'-H_2^0}_{pp\to ZH_2^0} = -\frac{g^4\alpha_s^2\Lambda_{ZZ'H_2^0}(v_\rho^2-v_\eta^2)\Pi\beta_{H_i^0}}{4096\pi^3\hat{s}^2c_W^2v_W^3}\mathrm{Re}\left[\frac{\chi^{(1)}(\hat{s})}{(s-m_{Z'}^2+im_{Z'}\Gamma_{Z'})}\sum_{q=u,d}m_q\left(m_u\frac{v_\eta}{v_\rho}-m_d\frac{v_\rho}{v_\eta}\right)T_3^q(1+2\delta_qI_q)\sum_{q=u,d}I_q^*\right], \tag{25}$$

which in Eq. (20) are considered the contribution of the Z' boson, in Eq. (21) the contribution of interference of H_1^0 and H_2^0, in Eqs. (22) and (23) the contribution of H_1^0 and H_2^0, in Eq. (24) the contribution of interference of Z' and H_1^0, and in Eq. (25) the contribution of interference of Z' and H_2^0. All these contributions are to produce the ZH_2^0. The sum runs over all generations, T_3^q is the quark weak isospin [$T_3^{u(d)} = +(-)1/2$], and Re stands for the real part of the expression. The loop function $I_i \equiv I(\delta_i = m_i^2/\hat{s})$, is defined by

$$I_i \equiv I_i(\delta_i) = \int_0^1 \frac{dx}{x}\ln\left[1-\frac{(1-x)x}{\delta_i}\right] = \begin{cases} -2\left[\sin^{-1}\left(\frac{1}{2\sqrt{\delta_i}}\right)\right]^2, & \delta_i > \frac{1}{4} \\ \frac{1}{2}\ln^2\left(\frac{r_+}{r_-}\right) - \frac{\pi^2}{2} + i\pi\ln\left(\frac{r_+}{r_-}\right), & \delta_i < \frac{1}{4}, \end{cases}$$

with $r_\pm = 1 \pm (1-4\delta_i)^{1/2}$ and $\delta_i = m_i^2/\hat{s}$. Here, $i = q$ stands for the particle (quark) running in the loop.

We have also defined Δ, Ω, and Π, which are equal to

$$\Delta = 4\hat{s} - \frac{\hat{u}^2}{2m_Z^2} + \frac{\hat{t}\,\hat{u}}{m_Z^2} - \frac{\hat{t}^2}{2m_Z^2}$$

$$\Omega = \frac{\hat{s}^2}{4m_Z^2} - \frac{\hat{s}}{2} + \frac{m_Z^2}{4} - \frac{\hat{s}m_{H_1^0}^2}{2m_Z^2} - \frac{m_{H_1^0}^2}{2} + \frac{m_{H_1^0}^4}{4m_Z^2}$$

$$\Pi = -\frac{\hat{s}^4}{8m_Z^4} + \frac{\hat{s}^3}{4m_Z^2} - \frac{\hat{s}^2\hat{u}}{8m_Z^2} - \frac{\hat{s}^2\hat{t}}{8m_Z^2} + \frac{\hat{s}^2}{8} + \frac{3\hat{s}\,\hat{u}}{8} + \frac{3\hat{s}\,\hat{t}}{8} - \frac{\hat{s}m_Z^2}{4} + \frac{\hat{s}^3m_{H_i^0}^2}{4m_Z^4} + \frac{\hat{s}^2m_{H_i^0}^2}{4m_Z^2} + \frac{\hat{s}\,\hat{u}\,m_{H_i^0}^2}{8m_Z^2} + \frac{\hat{s}\,\hat{t}\,m_{H_i^0}^2}{8m_Z^2} - \frac{3\hat{s}m_{H_i^0}^2}{4} - \frac{\hat{s}^2m_{H_i^0}^4}{4m_Z^4}.$$

The total cross section for the process $pp \to gg \to ZH_2^0$ is related to the subprocess $gg \to ZH_2^0$ total cross section $\hat{\sigma}$ through

$$\sigma = \int_{\tau_{\min}}^1 \int_{\ln\sqrt{\tau_{\min}}}^{-\ln\sqrt{\tau_{\min}}} d\tau dy G(\sqrt{\tau}e^y, Q^2)G(\sqrt{\tau}e^{-y}, Q^2)\hat{\sigma}(\tau, s), \tag{26}$$

where $G(x, Q^2)$ is the gluon structure function and $\tau_{\min}$ is given above.

IV. RESULTS AND CONCLUSIONS

In this work, we have calculated contributions regarding to the Drell–Yan and gluon-gluon fusion in the 3-3-1 model. We present the cross section for the process $pp \to ZH_2^0$ involving the Drell–Yan mechanism and the gluon-gluon fusion to produce such Higgs bosons for the LHC. In all calculations we take for the parameters and the VEV the following representative values, [12,17]: $\lambda_1 = 0.3078$, $\lambda_2 = 1.0$, $\lambda_3 = -0.025$, $\lambda_4 = 1.388$, $\lambda_5 = -1.567$, $\lambda_6 = 1.0$, $\lambda_7 = -2.0$, $\lambda_8 = -0.45$, $v_\eta = 195$ GeV, and for $\lambda_9 = -0.90(-0.76, -0.71)$, correspond $v_\chi = 1000(1500, 2000)$ GeV, these parameters are used to estimate the values for the particles masses which are given in Table I, it is to notice that the value of λ_9 was chosen this way in order to guarantee the approximation $-f \simeq v_\chi$ [12,17].

A. Higgs H_2^0

The Higgs H_2^0 in the 3-3-1 model is not coupled to a pair of standard bosons. It couples to quarks; leptons; $Z\,Z'$, $Z'Z'$ gauge bosons; $H_1^-H_1^+$, $H_2^-H_2^+$, h^0h^0, $H_1^0H_3^0$ Higgs bosons; V^-V^+ charged bosons; $U^{--}U^{++}$ double charged bosons; H_1^0Z, H_1^0Z' bosons; and $H^{--}H^{++}$ double-charged Higgs bosons [14]. The Higgs H_2^0 can be much heavier than 1017.2 GeV for $v_\chi = 1000$ GeV, 1525.8 GeV for $v_\chi = 1500$ GeV, and 2034.37 GeV for $v_\chi = 2000$ GeV, so the Higgs H_2^0 is a heavy particle. The coupling of the H_2^0 with H_1^0 contributes to the enhancement of the total cross section via the Drell–Yan and gluon-gluon fusion.

In Figs. 3 and 4, we show the cross section $pp \to ZH_2^0$ for Drell–Yan and gluon-gluon fusion; these processes will

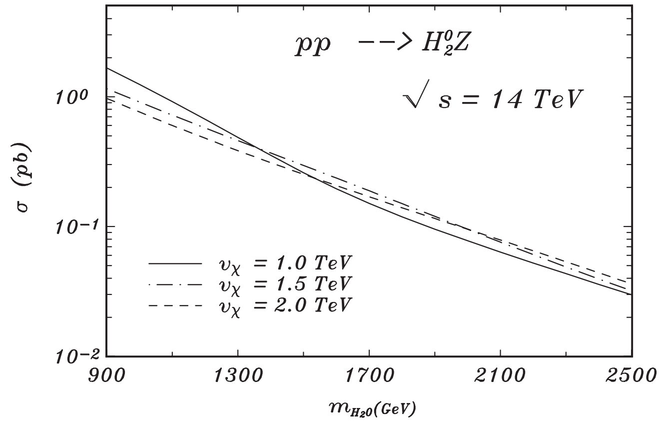

FIG. 3. Total cross section for the process $pp \to ZH_2^0$ as a function of $m_{H_2^0}$ at $\sqrt{s} = 14$ TeV via Drell–Yan. The solid line represents $v_\chi = 1.0$ TeV, the dotted-dashed line represents $v_\chi = 1.5$ TeV, and the short dashed line represents $v_\chi = 2.0$ TeV.

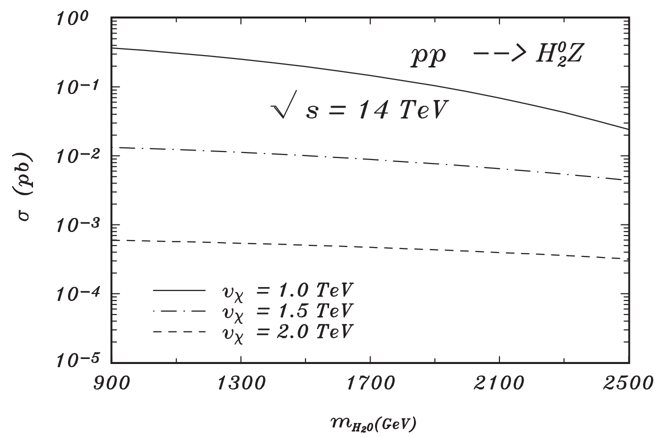

FIG. 4. Total cross section for the process $p\ p \to ZH_2^0$ as a function of $m_{H_2^0}$ at $\sqrt{s} = 14$ TeV via gluon-gluon fusion. The solid line represents $v_\chi = 1.0$ TeV, the dotted-dashed line represents $v_\chi = 1.5$ TeV, and the short dashed line represents $v_\chi = 2.0$ TeV.

be studied for $\sqrt{s} = 14$ TeV and for the vacuum expectation values $v_\chi = 1000$ GeV, $v_\chi = 1500$ GeV, and $v_\chi = 2000$ GeV. Considering that the expected integrated luminosity for the LHC collider that will be reached is of order of 300 fb^{-1}, then the statistics for $v_\chi = 1000$ GeV give a total of $\simeq 2.7 \times 10^5 (1.4 \times 10^5)$ events per year for Drell–Yan and $\simeq 9.3 \times 10^4 (7.5 \times 10^4)$ events per year for gluon-gluon fusion if we take the mass of the Higgs boson $m_{H_2^0} = 1100(1300)$ GeV ($\Gamma_{H_2^0} = 878.25,\ 1091.33$ GeV) and it corresponds to 14 TeV for the LHC, respectively. These values are in accord with Table I. It must be noticed that one must take care with large Higgs masses, as the width approaches the value of the mass itself for a very heavy Higgs, and one looses the concept of resonance.

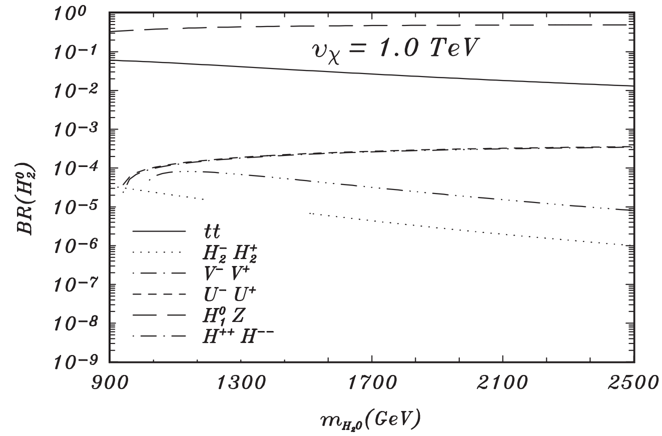

FIG. 5. Branching ratios (BRs) for the H_2^0 decays as functions of $m_{H_2^0}$ for $v_\chi = 1.0$ TeV.

To obtain event rates, we multiply the production cross sections by the respective branching ratios. Considering that the signal for $H_2^0 Z$ production for $m_{H_2^0} = 1100(1300)$ GeV and $v_\chi = 1000$ GeV will be $H_2^0 Z \to ZH_1^0 Z$, and taking into account that the branching ratios for these particles would be $BR(H_2^0 \to ZH_1^0) = 39.5(43.4)\%$ (see Fig. 5) and $BR(Z \to b\bar{b}) = 15.2\%$ and that the particles H_1^0 decay into W^+W^-, and taking into account that the branching ratios for these particles would be $BR(H_1^0 \to W^+W^-) = 23.1\%$ followed by leptonic decay of the boson W^+ into $\ell^+\nu$ and W^- into $\ell^-\bar{\nu}$, for which branching ratios for these particles would be $\mathrm{BR}(W \to \ell\nu) = 10.8\%$, then we would have approximately $\simeq 7(4)$ events per year for Drell–Yan and $\simeq 2(2)$ for gluon-gluon fusion for the signal $b\bar{b}b\bar{b}\ell^+\ell^-X$.

The statistics for $v_\chi = 1500$ give a total of $\simeq 7.1 \times 10^4 (4.5 \times 10^4)$ events per year for Drell–Yan and $\simeq 2.8 \times 10^3 (2.5 \times 10^3)$ events per year for gluon-gluon fusion if we take the mass of the Higgs boson $m_{H_2^0} = 1600(1800)$ GeV. These values are in accord with Table I. Taking into account the same signal as above, that is, $H_2^0 Z \to ZH_1^0 Z$, and taking into account that the branching ratios for these particles would be $BR(H_2^0 \to ZH_1^0) = 44.2(45.9)\%$ (see Fig. 6), $BR(Z \to b\bar{b}) = 15.2\%$, $BR(H_1^0 \to W^+W^-) = 23.1\%$, and $\mathrm{BR}(W \to \ell\nu) = 10.8\%$, we would have approximately $\simeq 2(1)$ events per year for Drell–Yan and $\simeq 0(0)$ for gluon-gluon fusion for the signal $b\bar{b}b\bar{b}\ell^+\ell^-X$.

With respect to vacuum expectation value $v_\chi = 2000$ GeV, for the masses of $m_{H_2^0} = 2100(2300)$, it will give a total of $\simeq 2.4 \times 10^4 (1.6 \times 10^4)$ events per year to produce H_2^0 for Drell–Yan, and with respect to gluon-gluon fusion, we will have $\simeq 119(107)$ events per year to produce the same particles. Taking into account the same signal as above, that is, $b\bar{b}b\bar{b}\ell^+\ell^-X$, and considering that the branching ratios for H_2^0 would be $BR(H_2^0 \to ZH_1^0) = 46.4(47.3)\%$ (see Fig. 7), we will have approximately

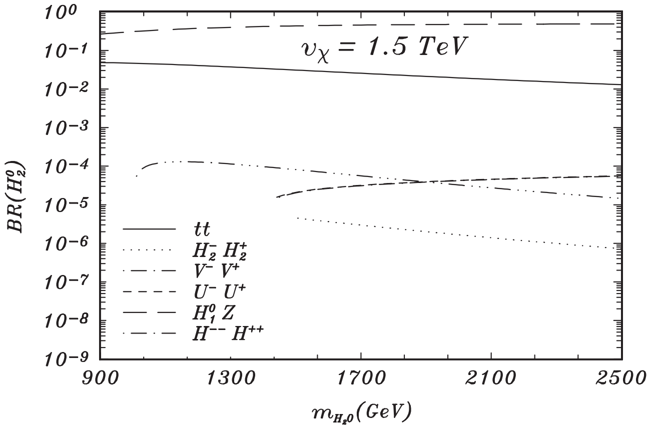

FIG. 6. BRs for the H_2^0 decays as functions of $m_{H_2^0}$ for $\upsilon_\chi = 1.5$ TeV.

$\simeq 0(0)$ events per year for Drell–Yan and $\simeq 0(0)$ for gluon-gluon fusion.

The main background to this signal is $t\bar{t}Z \to b\bar{b}b\bar{b}\ell^+\ell^- X$, for which the cross section at LO is $\simeq 1$ pb for $\sqrt{s} = 14$ TeV. Considering that the $t\bar{t}$ particles decay into $b\bar{b}W^+W^-$, for which the branching ratios for these particles would be $BR(t \to bW) = 99.8\%$ followed by leptonic decay of the boson W, that is, $BR(W \to \ell\nu) = 10.8\%$ and $BR(Z \to b\bar{b}) = 15.2\%$, then we would have approximately a total of $\simeq 530$ events for the background and $\simeq 9(6)$ events for the signal for $m_{H_2^0} = 1100(1300)$ GeV and $\upsilon_\chi = 1000$; on the other hand, for $\upsilon_\chi = 1500$ and $\upsilon_\chi = 2000$, the number of events for the signal is insignificant.

Therefore, we have that the statistical significance is $\simeq 0.39(0.26)\sigma$ for $m_{H_2^0} = 1100(1300)$ GeV, which is a low probability to detect the signals. The improvement will be significant if we consider a luminosity $\simeq 10$ times higher than original LHC design, which is what we are awaiting to happen for 2025; then, we will have $\simeq 90(60)$ events for the signals for $m_{H_2^0} = 1100(1300)$ GeV and $\upsilon_\chi = 1000$, which corresponds to having a $\simeq 3.9(2.6)\sigma$. Then, we have evidence for $\simeq 3.9\sigma$ discovery in the $b\bar{b}b\bar{b}\ell^+\ell^- X$ final state. On the other hand, for $\upsilon_\chi = 1500$, we have $\simeq 20(10)$ events for $m_{H_2^0} = 1600(1800)$ GeV, and this corresponds to $\simeq 0.89(0.44)\sigma$; for this last scenario, the signal is too small to be observed even with 3000 fb^{-1}.

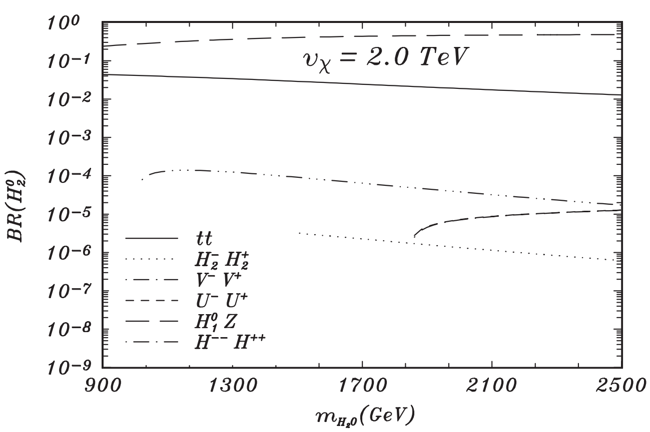

FIG. 7. BRs for the H_2^0 decays as functions of $m_{H_2^0}$ for $\upsilon_\chi = 2.0$ TeV.

To extract the signal from the background, we must select the $b\bar{b}$ channel using the techniques of b-flavor identification, thus reducing the huge QCD backgrounds of quark and gluon jets. Later, the Z that comes together with the H_2^0 and the other Z that comes from the decay of H_2^0 would appear as a peak in the invariant mass distribution of b-quark pairs. The charged lepton track from the W decay and the cut on the missing transverse momentum $\not{p}_T > 20$ GeV allows for a very strong reduction of the backgrounds.

The $H_2^0 Z$ will also decay into $t\bar{t}\,\ell^+\ell^-$, and considering that the branching ratios for these particles would be $BR(H_2^0 \to t\bar{t}) = 5.1(4.1)\%$ (see Fig. 5) and $BR(Z \to \ell^+\ell^-) = 3.4\%$ for the mass of the Higgs boson $m_{H_2^0} = 1100(1300)$ GeV and $\upsilon_\chi = 1000$ GeV and that the particles $t\bar{t}$ decay into $b\bar{b}W^+W^-$, for which the branching ratios for these particles would be $BR(t \to bW) = 99.8\%$, followed by leptonic decay of the boson W, that is, $BR(W \to e\nu) = 10.75\%$, then we would have approximately $\simeq 5(2)$ events per year for Drell–Yan. Regarding gluon-gluon fusion, we will have $\simeq 2(1)$ events per year to produce the same particles. Considering the vacuum expectation value $\upsilon_\chi = 1500$ GeV and the branching ratios $BR(H_2^0 \to t\bar{t}) = 2.8(2.3)\%$ (see Fig. 6) and taking the same parameters and branching ratios for the same particles given above, then we would have for $m_{H_2^0} = 1600(1800)$ a total of $\simeq 1(1)$ events of H_2^0 produced per year for Drell–Yan, and in respect to gluon-gluon fusion, the number of events per year for the signal will be $\simeq 0(0)$.

Taking again the irreducible background $t\bar{t}Z \to b\bar{b}e^+e^-\ell^+\ell^- X$ and using CompHep [18], we have that a cross section at LO is $\simeq 1$ pb, which gives $\simeq 117$ events. So, we will have a total of $\simeq 6(3)$ events per year for the signal for $m_{H_2^0} = 1100(1300)$ GeV and $\upsilon_\chi = 1000$, and for $\upsilon_\chi = 1500$ and $\upsilon_\chi = 2000$, the number of events is insignificant.

Then, we have that the statistical significance is $\simeq 0.55(0.28)\sigma$ for $m_{H_2^0} = 1100(1300)$ GeV and $\upsilon_\chi = 1000$ GeV. For this scenario, the signal significance is smaller than 1σ, and discovery cannot be accomplished unless the luminosity will be improved. So, if we enhance the integrated luminosity up to 3000 fb^{-1}, then we will have $\simeq 60(30)$ events for the signals for $m_{H_2^0} = 1100(1300)$ GeV and $\upsilon_\chi = 1000$, which corresponds to having a $\simeq 5.5(2.8)\sigma$ discovery in the $b\bar{b}e^+e^-\ell^+\ell^- X$ final state, and for $\upsilon_\chi = 1500(2000)$, the signal will not be visible in this channel. We impose the following cuts to improve the statistical significance of a signal; i.e., we isolate a hard lepton from the W decay with $p_T^\ell > 20$ GeV, put the cut on the missing transverse momentum $\not{p}_T > 20$ GeV, and apply the Z window cut $|m_{\ell^+\ell^-} - m_Z| > 10$ GeV, which removes events for which the leptons come from Z decay [19]. However, all these scenarios can only be cleared by a careful Monte Carlo work to determine the size of the signal and background.

In summary, we showed in this work that, in the context of the 3-3-1 model, the signatures for neutral Higgs boson H_2^0 can be significant in the LHC collider if we take $v_\chi = 1000$, $m_{H_2^0} = 1100(1300)$ GeV, $\sqrt{s} = 14$ TeV, and a luminosity of 3000 fb^{-1}. In other scenarios, the signal will be too small to be observed even with 3000 fb^{-1}. Our study indicates the possibility of obtaining a signal of this new particle in the channel $t\bar{t}Z \to b\bar{b}e^+e^-\ell^+\ell^-X$. If this model is realizable in the nature, certainly new particles will appear such as H_2^0, Z' in the context of this study.

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