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CONSTRAINTS ON THE $SU(3)_C \otimes SU(2)_L \otimes SU(3)_R \otimes U(1)_X$ MODEL FROM THE PROCESS $W_2 \to t\bar{b}$

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Abstract. *We study the process $pp \to W_2 \to t\bar{b}$ in the $SU(3)_C \otimes SU(2)_L \otimes SU(3)_R \otimes U(1)_X$ model. The production cross-section times branching ratio is determined. We use the LHC data to derive constraints on the new physical scale, m_{W_2}, and the elements of the right-handed quark mixing matrix.*

Keywords: 3-2-3-1 model, standard model.

Classification numbers: 12.15.24, 11.30.Qx, 12.15.Ff.

I. INTRODUCTION

The minimal left-right symmetric theory of the weak interactions is attractive extensions of the standard model (SM), which provides natural explanations for electroweak parity asymmetry and small neutrino masses [1]. The appearance of flavor changing neutral currents allows explaining the neutral meson mixing and rare meson decay [2]. The new physics provides an answer to the V_{ub} problem [3]. However, the minimal left-right model does not provide natural solutions for dark matter and family number. Actually, the lightest right-handed neutrino is considered as long live warm dark matter. Because of the gauge interactions, the lightest right-handed neutrino would overpopulate the universe [4]. Besides, the tableau of cold dark matter has also been investigated by adding a new field or imposing a new symmetry, the global and Abelian symmetries of the minimal left-right symmetric model [5]. Even so, the dark matter sector decouples from the normal sector due to gauge symmetry. It means that the dark matter still is arbitrary and ad hoc.

Contrary to previous expansions, we found a stabilizing mechanism of the dark matter via considering a non-cummutative $B-L$ gauge symmetry that warrants the dark and normal matters are defined in the same irreducible gauge multiplet [6]. But the dark matter has a wrong $B-L$ number and the normal matter has a normal $B-L$ number. After the $B-L$ symmetry breaking, only the dark matter transforms non-trivially under a residual discrete gauge symmetry (W-parity) that protects the stability of dark matter. Parallel to this, the vacuum expectation value (VEV) which breaks $B-L$ symmetry, also defines a seesaw scale. In order to create a non-cummutative $B-L$ operator, we can expand non-commutative symmetry by enlarging the left, right weak-isospin groups [6–9]. Following this research direction, both dark matter and family number problems can be solved simultaneously.

Let us consider the simplest theory, which is given by the $SU(3)_C \times SU(2)_L \times SU(3)_R \times U(1)_X$(3-2-3-1) gauge symmetry. This is called the simplest theory because the left-handed fermion content and symmetry are the same as those of the standard model and minimal left-right symmetric model, but the right sector is extended. Thus, the 3-2-3-1 model can explain the family number problem via considering a $SU(3)_R$ anomaly cancellation. The small neutrino mass can be solved by seesaw mechanism. The highlight of the 3-2-3-1 model is that the $B-L$ charge is determined as: $\frac{B-L}{2} = \beta T_{8R} + X$. It behaves as the non-commuting gauge charge of $SU(3)_R \times U(1)_X$. The VEV, which breaks the $B-L$ symmetry, defines both a seesaw scale and $W-$ parity. The particles with non-trivial $B-L$ quantum number carry a W-odd parity and the remaining particles carry a $W-$ even parity. The lightest $W-$odd particle is identified to the dark matter. The flavor changing neutral currents (FCNCs) in the 3-2-3-1 model arises at the tree level via the interactions of neutral new gauge boson and Higgs boson which are powerful probes of physics beyond the standard model such as: the $K^0-\bar{K}^0$, $B-\bar{B}$ masses difference at the unacceptable and other new physics [8].

Recently, the ATLAS and CMS collaboration report excesses in searches for new gauge bosons, W', Z' via various decay modes. The new gauge bosons decay into a pair of leptons [10,11], diboson [12], dijet [13]. Several studies on the Z_R, Z'_R production and their decay channels in the 3-2-3-1 model have been studied in [8]. The results exclude the Z_R, Z'_R bosons lighter than 4.0 TeV. Notably, search for $W' \to t\bar{b}$ decays in the hadronic final state using pp collision at $\sqrt{s} = 13$ TeV with Atlas detector is presented [14]. No clear signal for this decay mode has been found and Atlas collaboration provides upper limits on production cross sections times branching ratio as a function of unknown W' mass. The 3-2-3-1 model also predicts the existence of the W_2 gauge bosons. It is the mediator of a new charged vector current that can be heavy enough to decay into a t-quark and a $b-$quark. In this work, we discuss how these limits can be used to constrain the W_2 in the 3-2-3-1 model.

The rest of this paper is organized as follows: In Sec. **II**, we give a review of the model. Section **III** studies the W_R- boson production and decay. The total cross section times branching ratio is considered as a function of the W_2 mass and the elements of the right-handed quark mixing matrix. We conclude this work in Sec. **IV**.

II. THE MODEL

As state above, the 3-2-3-1 model relies on the gauge group $SU(3)_C \otimes SU(2)_L \otimes SU(3)_R \otimes U(1)_X$. Compared to the standard model, the hypercharge is extended as $U(1)_Y \to SU(3)_R \times U(1)_X$. This extension allows resolving the small neutrino mass, dark matter, and family number problems. The electric charge and hypercharge are embedded in the gauge symmetry as follows $Q = T_{3L} + T_{3R} + \beta T_{8R} + X, \quad Y = T_{3R} + \beta T_{8R} + X$, where T_{aL} ($a = 1,2,3$), T_{iR} ($i = 1,2,3,...,8$), and X are $SU(2)_L$, $SU(3)_R$, and $U(1)_X$ generators, respectively. The coefficient β can be expressed via an electric charge parameter (q) as $\beta = -(2q+1)/\sqrt{3}$. The quarks and fermions transform under the gauge group as $\Psi_{aL} = (\nu_{aL}, e_{aL})^T \sim (1,2,1,-\frac{1}{2})$, $Q_{aL} = (u_{aL}, d_{aL})^T \sim (3,2,1,\frac{1}{6})$ and $\Psi_{aR} = (\nu_{aR}, e_{aR}, E^q_{aR})^T \sim (1,1,3,\frac{q-1}{3})$, $Q_{\alpha R} = (d_{\alpha R}, -u_{\alpha R}, J^{-q-\frac{1}{3}}_{\alpha R})^T \sim (3,2,1,\frac{1}{6})$, $Q_{3R} = (u_{3R}, d_{3R}, J^{q+\frac{2}{3}}_{3R})^T \sim (3,1,3,\frac{q+1}{3})$, $E^q_{aL} \sim (1,1,1,q)$, $J^{q+\frac{2}{3}}_{3L} \sim (3,1,1,q+\frac{2}{3})$, $J^{-q-\frac{1}{3}}_{\alpha L} \sim (3,1,1,-q-\frac{1}{3})$, where $a = 1,2,3$ and $\alpha = 1,2$ are family indices. ν_R, E and J are new particles included to complete the representations. The minimal Higgs content, which breaks the gauge symmetry and generates the particle masses appropriately, consists a bi-triplet $S \sim (1,2,3^*,-\frac{2q+1}{6})$, a sextet $\Xi \sim (1,1,6,\frac{2(q-1)}{3})$ and a triplet $\phi \sim (1,1,3,-\frac{2q+1}{3})$. Spontaneous breaking of gauge symmetry $SU(2)_L \times SU(3)_R \times U(1)_X$ down to $U(1)_{EM}$ is achieved by vacuum expectation values (VEVs) of the neutral fields that are denoted by

$$\langle S \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} u & 0 & 0 \\ 0 & v & 0 \end{pmatrix}, \quad \langle \phi \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ 0 \\ w \end{pmatrix}, \quad \langle \Xi \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} \Lambda & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 0 \end{pmatrix}. \tag{1}$$

Among the VEVs u, v and w, Λ, the hierarchy $u, v \ll w, \Lambda$ is necessary to protect the success of the standard model. The full Higgs potential and the scalar mass spectrum are studied in [8]. The Yukawa Lagrangian involving the quark fields is given by

$$\mathcal{L}^{quark}_{\text{Yukawa}} = h^q_{a3} \bar{Q}_{aL} S Q_{3R} + h^q_{a\beta} \bar{\tilde{Q}}_{aL} S^* Q_{\beta R} + h^J_{33} \bar{J}_{3L} \phi^\dagger Q_{3R} + h^J_{\alpha\beta} \bar{J}_{\alpha L} \phi^T Q_{\beta R} + H.c. \tag{2}$$

The first two terms lead to the following mass matrices for the up-type and down-type quarks

$$\mathcal{M}^u = -\frac{1}{2} \begin{pmatrix} h^q_{11} v & h^q_{12} v & h^q_{13} u \\ h^q_{21} v & h^q_{22} v & h^q_{23} u \\ h^q_{31} v & h^q_{32} v & h^q_{33} u \end{pmatrix}, \quad \mathcal{M}^d = -\frac{1}{2} \begin{pmatrix} h^q_{11} u & h^q_{12} u & h^q_{13} v \\ h^q_{21} u & h^q_{22} u & h^q_{23} v \\ h^q_{31} u & h^q_{32} u & h^q_{33} v \end{pmatrix}. \tag{3}$$

The matrices can be diagonalized by using bi-unitary transformations as follows

$$V^\dagger_{dL} \mathcal{M}^d V_{dR} = M^d = Diag(m_d, m_s, m_b), \quad V^\dagger_{uL} \mathcal{M}^u V_{uR} = M^u = Diag(m_u, m_c, m_t). \tag{4}$$

In the charged gauge boson sector, the left and right-handed W bosons mix via a 2×2 matrix. The physical states are presented as follows

$$W^\pm_{1\mu} = c_\xi W^\pm_{L\mu} - s_\xi W^\pm_{R\mu}, \quad W^\pm_{2\mu} = s_\xi W^\pm_{L\mu} + c_\xi W^\pm_{R\mu}, \tag{5}$$

where $c_\xi = \cos\xi, s_\xi = \sin\xi$ and the mixing angle ξ is defined by $t_{2\xi} = \tan 2\xi = \frac{-4t_R uv}{2t_R^2\Lambda^2 + (t_R^2-1)(u^2+v^2)}$ with $t_R = \frac{g_R}{g_L}$. The mass eigenvalues of the $W_1^\pm, W_2^\pm$ gauge bosons are given by

$$\begin{aligned} m_{W_1}^2 &\simeq \frac{g_L^2}{4}\left[u^2+v^2-\frac{4t_R^2u^2v^2}{2t_R^2\Lambda^2+(t_R^2-1)(u^2+v^2)}\right], \\ m_{W_2}^2 &\simeq \frac{g_R^2}{4}\left[u^2+v^2+2\Lambda^2+\frac{4t_R^2u^2v^2}{2t_R^2\Lambda^2+(t_R^2-1)(u^2+v^2)}\right]. \end{aligned} \tag{6}$$

Because of the condition $u, v \ll w, \Lambda$, the W_1 boson has a small mass in the weak scales (u, v) which is identified to the standard model W boson, whereas the W_2 boson is a new, heavy charged gauge boson with mass proportional to Λ scale. The mixing between these two fields is small since $\xi \to 0$ due to the above condition. The coupling of $W_{2\mu}^\pm$ to quarks (mass eigenstates) is written by

$$\frac{W_{2\mu}^+}{\sqrt{2}}\bar{u}'\left(g_R c_\xi V_{CKM}^R P_R + g_L s_\xi V_{CKM}^L P_L\right)\gamma^\mu d' + H.c., \tag{7}$$

where $V_{CKM}^L = V_{uL}^\dagger V_{dL}, V_{CKM}^R = V_{uR}^\dagger V_{dR}$ and $u' = (u,c,t)^T$, $d' = (d,s,b)^T$. Similar to the quark part, the Yukawa Lagrangian involving the lepton fields is given as

$$\mathscr{L}_{\text{Yukawa}}^{lepton} = h_{ab}^l \bar{\Psi}_{aL} S \Psi_{bR} + h_{ab}^R \bar{\Psi}_{aR}^c \Xi^\dagger \Psi_{bR} + h_{ab}^E \bar{E}_{aL}\phi^\dagger \Psi_{bR} + H.c.. \tag{8}$$

The Dirac neutrino mass matrix $[m_D]_{ab} = -h_{ab}^l \frac{u}{\sqrt{2}}$ and the Majorana neutrino mass matrix $[M_R]_{ab} = -\sqrt{2}h_{ab}^R\Lambda$. The active neutrinos have a small mass via a type I seesaw mechanism. The weak states for the left and right-handed neutrinos respectively are denoted as $\nu_L = (\nu_{eL}, \nu_{\mu L}, \nu_{\tau L})^T, \nu_R = (\nu_{eR}, \nu_{\mu R}, \nu_{\tau R})^T$. The light neutrino mass eigenstates $\nu = (\nu_{1L}, \nu_{2L}, \nu_{3L})^T$ and heavy neutrino mass eigenstate $N = (N_{1R}, N_{2R}, N_{3R})^T$ which relate to the weak states as: $\nu_L = U_L\nu, \nu_R = U_R N$. If the charged leptons by themselves are physical fields, the interaction of new charged gauge boson, W_2, with the normal leptons is given as follows

$$\frac{W_{2\mu}^+}{\sqrt{2}}\bar{N}\left(g_R c_\xi U_R P_R\right)\gamma^\mu l + \frac{W_{2\mu}^+}{\sqrt{2}}\left(g_L s_\xi U_L P_L\right)\gamma^\mu l + H.c.. \tag{9}$$

Thus the U_L matrix is identified to the Maki-Nakagawa-Sakata matrix.

III. SEARCH FOR $W_2 \to tb$ DECAY

At the LHC, the W_2 gauge boson is produced via the interactions given in Eq.(2). The production cross section at the LHC with the center of mass energy $\sqrt{s} = 13$ TeV is evaluated as follows

$$\sigma(pp \to W_2X) \simeq \frac{\pi}{48s}\sum_{ij}\left[\left(C_{q_{ij}}^L\right)^2 + \left(C_{q_{ij}}^R\right)^2\right] w_{ij}\left(\frac{M_{W_2}^2}{s}, M_{W_2}\right), \tag{10}$$

where $C_{q_{ij}}^L = g_L\left(V_{CKM}^L\right)s_\xi, C_{q_{ij}}^R = g_R\left(V_{CKM}^R\right)_{ij}c_\xi$ and the function $w_{ij}\left(\frac{M_{W_2}^2}{s}, M_{W_2}\right)$ is defined as

$$w_{ij}\left(\frac{M_{W_2}^2}{s}, M_{W_2}\right) = \int_z^1 \frac{dx}{x}\left[u_i(x,\mu)\bar{d}_j(\frac{z}{x},\mu) + \bar{u}_i(x,\mu)d_j(\frac{z}{x},\mu)\right]. \tag{11}$$

Denoting that $u_i(x,\mu)$ and $d_i(x,\mu)$ are a parton distribution inside the proton. Due to the interaction terms are given in Eqs. (2) and (9), the W_2 gauge boson decays into a pair of lepton, two quark jets. The W_2 gauge boson mixes with the W_1 gauge boson, thus it couples to the standard model Higgs (h) and neutral gauge boson Z. In the limit $m_Z < m_{W_2} < m_{N_i}$, the W_2 gauge boson can not decay into the right-handed neutrinos. The W_2 partial decay widths are approximately given by

$$\Gamma(W_2 \to q_i\bar{q}_j) = \frac{m_{W_2}}{16\pi}\left(\left|g_L s_\xi (V^L_{CKM})_{ij}\right|^2 + \left|g_L c_\xi (V^R_{CKM})_{ij}\right|^2\right), \tag{12}$$

$$\Gamma(W_2 \to t\bar{q}_i) = \frac{\left(1-\frac{m_t^2}{m_{W_2}^2}\right)^2}{16\pi m_{W_2}}\left(\left|g_L s_\xi (V^L_{CKM})_{3i}\right|^2 + \left|g_L c_\xi (V^R_{CKM})_{3i}\right|^2\right)\left(m_{W_2}^2 + \frac{m_t^2}{2}\right), \tag{13}$$

$$\Gamma(W_2 \to l\bar{\nu}) = \frac{m_{W_2}}{48\pi}\left(\left|g_L s_\xi (U_L)_{ij}\right|^2 + \left|g_L c_\xi (U_R)_{ij}\right|^2\right), \tag{14}$$

$$\Gamma(W_2 \to W_1 Z) \simeq \frac{m_{W_2}}{192\pi} g^2_{W_2W_1Z}\frac{m^4_{W_2}}{m^2_{W_1}m^2_Z}, \qquad \Gamma(W_2 \to W_1 h) \simeq \frac{m_{W_2}}{192\pi}\left(\frac{g_{W_2W_1h}}{m_{W_1}}\right)^2, \tag{15}$$

where $g_{W_2W_1Z} = \frac{s_{2\xi}}{2t_R\cos\theta_W}$, $g_{W_2W_1h} = \frac{g_L^2\left((1-t_R^2)(u^2+v^2)s_{2\xi}-4t_R uvc_{2\xi}\right)}{4\sqrt{u^2+v^2}}$. The θ_W is a Weinberg angle.

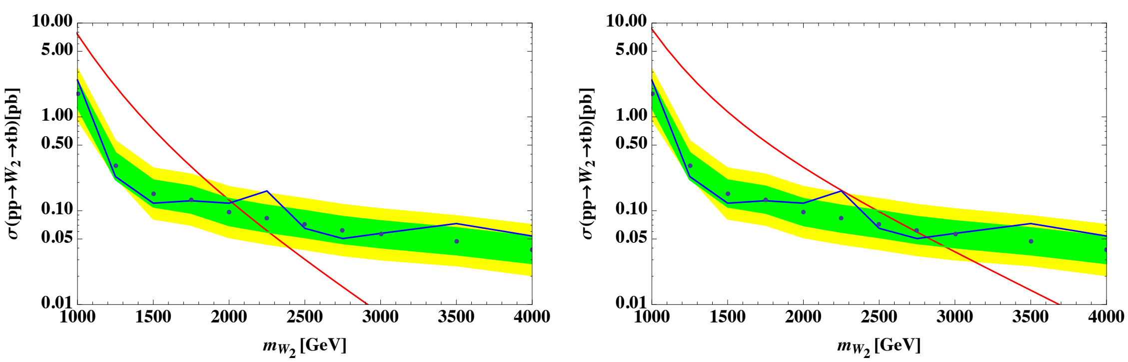

Fig. 1. The cross section, $\sigma(pp \to \mathscr{W}_2 \to t\bar{b})$, as a function of $m_{\mathscr{W}_2}$ with assuming $g_L = g_R, u = 220$ GeV,$\beta = \pm\frac{1}{\sqrt{3}}$. The yellow and green regions are the expected 95% CL limits with ± 1 and ± 2 standard deviation bands respectively. The dots present the expected 95% CL limit while the blue lines are the observed 95% CL limit [14]. The theoretical prediction for the models with $\pm\frac{1}{\sqrt{3}}$ is shown by red line.

In what follows, we consider the total cross section $\sigma(pp \to W_2).Br(W_2 \to t\bar{b})$ at the LHC for $\sqrt{s} = 13$ TeV and the parton distributions are obtained from [15]. Figure 1, we show the cross-sections times branching ratio of W_2 to $t\bar{b}$ decay for versions with $\beta = \pm\frac{1}{\sqrt{3}}$. The experimental search for $W_2 \to t\bar{b}$ uses $36.1 fb^{-1}$ of $\sqrt{s} = 13$ TeV proton-proton collision, which collected by

Atlas detector at the LHC, gives a negative signal [14]. In the case of fixed the $V^{L,R}_{CKM}$ matrices, it can be converted into the lower limit on the W_2 mass. The difference between the mass exclusion limit results for W_2 signals is due to differential choice of the V^R_{CKM} matrix. In particular, the left panel of the Figs.(1) is considered by taking $V^L_{CKM} = V_{CKM}, V^R_{CKM} = I$. The lower limit of W_2 mass is $m_{W_2} \simeq 2$ TeV. In the right panel, we take $V^L_{CKM} = V_{CKM}, [V^R_{CKM}]_{ij} = 0.999$ for $i = j$ and $[V^R_{CKM}]_{ij} = 0.001$ for the remaining matrix elements, we obtain the lower limit on the W_2 mass, $m_{W_2} \simeq 2.8$ TeV.

IV. CONCLUSION

We have examined the production and decay processes of W_2 gauge boson in the 3-2-3-1 model. Based on the interactions of W_2 gauge boson, we have given the analytical expressions for W_2 production cross section and partial decay widths. The W_2 gauge boson cross section times branching ratio of $t\bar{b}$ decay has been numerical studied. The predicted results depend on not only the W_2 gauge boson mass, but also on the mixing matrix of right-handed quarks. Compared to the collected data of Atlas detector at the LHC with proton-proton collision energy $\sqrt{s} = 13$ TeV, we can obtain the lower limit of the W_2 gauge boson mass. Because no significant deviation from the Standard Model prediction was observed, it leads to exclude W_2 boson with mass below a few TeV at the 95%. The exclusion limit for W_2 gauge boson mass depends on the choice of V^R_{CKM} matrix.

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