

Heavy Higgs bosons and the 2 TeV W' boson

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Abstract

The hints from the LHC for the existence of a W' boson of mass around 1.9 TeV point towards a certain $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge theory with an extended Higgs sector. We show that the decays of the W' boson into heavy Higgs bosons have sizable branching fractions. Interpreting the ATLAS excess events in the search for same-sign lepton pairs plus b jets as arising from W' cascade decays, we estimate that the masses of the heavy Higgs bosons are in the 400–700 GeV range.

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1 Introduction

Using LHC data at $\sqrt{s} = 8$ TeV, the ATLAS and CMS Collaborations have reported deviations from the Standard Model (SM) of statistical significance between 2σ and 3σ in several final states, indicating mass peaks in the 1.8–2 TeV range [1]–[5]. The cross sections required for producing these mass peaks are consistent with the properties of a W' boson in an $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge theory with right-handed neutrinos that have Dirac masses at the TeV scale [6].

The spontaneous breaking of $SU(2) \times SU(2) \times U(1)$ gauge groups requires an extended Higgs sector. For large regions of parameter space, the W' boson has large branching fractions into heavy scalars from the Higgs sector [7][8]. We show here that the W' boson hinted by the LHC data may also decay into $H^+ A^0$ and $H^+ H^0$, where H^+ , A^0 and H^0 are heavy spin-0 particles present in Two-Higgs-Doublet models. We compute the branching fractions for these decays and present some evidence that signals for the $W' \rightarrow H^+ A^0/H^0$ processes may already be visible in the 8 TeV LHC data.

There are numerous and diverse studies of $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ models, spanning four decades [9]. An interesting aspect of the left-right symmetric models is that they can be embedded in the minimal $SO(10)$ grand unified theory. This scenario must be significantly modified due to the presence of Dirac masses for right-handed neutrinos required by the CMS e^+e^-jj events.¹ The theory introduced in [6] involves at least one vectorlike fermion transforming as a doublet under $SU(2)_R$. This may be part of an additional $SO(10)$ multiplet, but it may also be associated with entirely different UV completions.

In Section 2 we write down the Higgs potential and analyze its implications for the scalar spectrum. In Section 3 we derive the interactions of the W' boson with scalars, and compute all W' branching fractions. The couplings of heavy Higgs bosons to quarks are discussed in Section 4. LHC signals of heavy scalars produced in W' decays are the subject of Section 5. Our conclusions are summarized in Section 6.

¹Unless the right-handed neutrinos have TeV-scale masses with the split between two of them at the MeV scale [10].

2 Extended Higgs sector

The Higgs sector of the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge theory discussed in [6] consists of two complex scalar fields: an $SU(2)_R$ triplet T of $B - L$ charge $+2$, and an $SU(2)_L \times SU(2)_R$ bidoublet Σ of $B - L$ charge 0 .

2.1 Scalar potential

The renormalizable Higgs potential is given by

$$V(T) + V(T, \Sigma) + V(\Sigma) \quad , \quad (2.1)$$

where the triplet-only potential is

$$V(T) = -M_T^2 \text{Tr}(T^\dagger T) + \frac{\lambda_T}{2} [\text{Tr}(T^\dagger T)]^2 + \frac{\lambda'_T}{2} \text{Tr}[(T^\dagger T)^2] \quad . \quad (2.2)$$

The bidoublet-only potential $V(\Sigma)$ is chosen such that by itself it does not generate a VEV for Σ . This is discussed later, together with $V(T, \Sigma)$, which collects all the terms that involve both scalars.

For $M_T^2 > 0$ and $\lambda_T + \lambda'_T > 0$, the T scalar acquires a VEV, which upon an $SU(2)_R \times U(1)_{B-L}$ transformation can be written as

$$\langle T \rangle = \begin{pmatrix} 0 & 0 \\ u_T & 0 \end{pmatrix} \quad , \quad (2.3)$$

where the minimization of $V(T)$ gives

$$u_T = \frac{M_T}{\sqrt{\lambda_T + \lambda'_T}} > 0 \quad . \quad (2.4)$$

This breaks $SU(2)_R \times U(1)_{B-L}$ down to the SM hypercharge gauge group, $U(1)_Y$, leading to large masses for the W' and Z' bosons. The value of the T VEV is related to the parameters of the W' boson. In the next section we will show that the parameters indicated by the LHC mass peaks near 2 TeV imply $u_T \approx 3-4$ TeV.

The triplet field includes 6 degrees of freedom, and can be written as $T = (T_1, T_2, T_3)$, with T_i ($i = 1, 2, 3$) complex scalars, or more explicitly as

$$T = \sum_{i=1}^3 T_i \sigma_i \equiv \begin{pmatrix} \frac{1}{\sqrt{2}} G_R^+ & T^{++} \\ u_T + \frac{1}{\sqrt{2}} (T^0 + i G_R^0) & -\frac{1}{\sqrt{2}} G_R^+ \end{pmatrix} \quad , \quad (2.5)$$

where σ_i are the Pauli matrices, and the factors of $\sqrt{2}$ are required for canonical normalization of the kinetic terms. The fields of definite electric charge, which are combinations of the T_i components, include three Nambu-Goldstone bosons (G_R^\pm, G_R^0). These become the longitudinal degrees of freedom of the W'^\pm and Z' bosons. The three remaining fields are a real scalar T^0 , a doubly-charged scalar T^{++} and its charge conjugate state T^{--} . These have masses given by

$$M_{T^0} = \sqrt{\lambda_T + \lambda'_T} u_T \quad , \quad M_{T^{++}} = \sqrt{\lambda'_T} u_T \quad . \quad (2.6)$$

For quartic couplings in the 0.1–1 range and in the absence of fine tuning, the T^0 and T^{++} particles have masses comparable to, or heavier than W' .

The bidoublet-only potential includes the following terms:

$$\begin{aligned} V(\Sigma) = & M_\Sigma^2 \text{Tr}(\Sigma^\dagger \Sigma) + \frac{\lambda_\Sigma}{2} [\text{Tr}(\Sigma^\dagger \Sigma)]^2 + \frac{\tilde{\lambda}_\Sigma}{4} \left| \text{Tr}(\tilde{\Sigma}^\dagger \Sigma) \right|^2 \\ & + \left[\frac{1}{2} \tilde{M}_\Sigma^2 \text{Tr}(\tilde{\Sigma}^\dagger \Sigma) + \frac{\tilde{\lambda}'_\Sigma}{8} \left(\text{Tr}(\tilde{\Sigma}^\dagger \Sigma) \right)^2 + \frac{\tilde{\lambda}''_\Sigma}{2} \text{Tr}(\Sigma^\dagger \Sigma) \text{Tr}(\Sigma^\dagger \tilde{\Sigma}) + \text{H.c.} \right] \quad , \quad (2.7) \end{aligned}$$

where $\tilde{\Sigma}$ is the charge conjugate state of the bidoublet. Other terms, such as $\text{Tr}[(\Sigma^\dagger \Sigma)^2]$ can be written as linear combinations of the terms in Eq. (2.7). We take $M_\Sigma^2, \tilde{M}_\Sigma^2 > 0$, which is a sufficient condition for the minimum of $V(\Sigma)$ to be at $\langle \Sigma \rangle = 0$. The mixed terms, which involve both the T and Σ scalars, will induce a nonzero VEV. In terms of fields of definite electric charge, the bidoublet scalar can be written as

$$\Sigma = \begin{pmatrix} \Sigma_1^0 & \Sigma_2^+ \\ \Sigma_1^- & \Sigma_2^0 \end{pmatrix} \quad , \quad (2.8)$$

and its charge conjugate state is

$$\tilde{\Sigma} = \sigma_2 \Sigma^* \sigma_2 = \begin{pmatrix} \Sigma_2^{0*} & -\Sigma_1^+ \\ -\Sigma_2^- & \Sigma_1^{0*} \end{pmatrix} \quad . \quad (2.9)$$

The mixed terms in the potential are:

$$V(T, \Sigma) = -\lambda_{T\Sigma} \text{Tr}(T^\dagger T \Sigma^\dagger \Sigma) - \tilde{\lambda}_{T\Sigma} \text{Tr}(T^\dagger T \tilde{\Sigma}^\dagger \tilde{\Sigma}) - \left(\frac{\tilde{\lambda}'_{T\Sigma}}{2} \text{Tr}(T^\dagger T) \text{Tr}(\Sigma^\dagger \tilde{\Sigma}) + \text{H.c.} \right) \quad (2.10)$$

For $\lambda_{T\Sigma} > M_\Sigma^2/u_T^2$ the VEV of T induces a negative squared mass for (Σ_2^0, Σ_2^-) , given by $-\lambda_{T\Sigma} u_T^2 + M_\Sigma^2 < 0$. Likewise, the squared mass for (Σ_1^0, Σ_1^-) turns negative, due to the

second term in Eq. (2.10), when $-\tilde{\lambda}_{T\Sigma}u_T^2 + M_\Sigma^2 < 0$. In addition, a $\Sigma_1^0\Sigma_2^0$ term is induced for $\tilde{\lambda}'_{T\Sigma} \neq 0$. Thus, for a range of parameters, the VEV of Σ takes the form

$$\langle \Sigma \rangle = v_H \begin{pmatrix} \cos\beta & 0 \\ 0 & e^{i\alpha_\Sigma} \sin\beta \end{pmatrix}, \quad (2.11)$$

where $v_H \simeq 174$ GeV is the electroweak scale. We are interested in the case where $u_T/v_H \sim 20$. The effect of the Σ VEV on the T^0 and T^{++} masses and couplings is thus negligible. At energy scales below the T^0 and T^{++} masses, the scalar sector consists only of Σ , which is the same as two Higgs doublets.

2.2 Higgs bosons

Using the notation of Two-Higgs-Doublet models [11], the components of Σ defined in Eq. (2.8) are related to the Higgs doublets H_1 and H_2 as follows:

$$H_1 = \begin{pmatrix} -\Sigma_1^+ \\ \Sigma_1^{0*} \end{pmatrix} = \begin{pmatrix} -H^+ \sin\beta + G^+ \cos\beta \\ v_H \cos\beta + \frac{1}{\sqrt{2}}(-h^0 \sin\alpha + H^0 \cos\alpha - iA^0 \sin\beta + iG^0 \cos\beta) \end{pmatrix},$$

$$H_2 = \begin{pmatrix} \Sigma_2^+ \\ \Sigma_2^0 \end{pmatrix} = \begin{pmatrix} H^+ \cos\beta + G^+ \sin\beta \\ v_H e^{i\alpha_\Sigma} \sin\beta + \frac{1}{\sqrt{2}}(h^0 \cos\alpha + H^0 \sin\alpha + iA^0 \cos\beta + iG^0 \sin\beta) \end{pmatrix}, \quad (2.12)$$

where h^0 is the SM-like Higgs boson, and G^\pm, G^0 are the Nambu-Goldstone bosons, which become the longitudinal W and Z . We have not included here the effects of the CP-violating phase α_Σ , which would lead to $H^0 - A^0$ mass mixing. The measurements of h^0 couplings are in good agreement with the SM predictions, implying an alignment limit [12], $\alpha = \beta - \pi/2$. In addition, the observed ATLAS events [1] consistent with the $W' \rightarrow WZ$ process indicate $\sin 2\beta \gtrsim 0.8$ (see Ref. [6]), which gives $0.5 \lesssim \tan\beta \lesssim 2$.

We have obtained a Two-Higgs-Doublet model with a potential formed of the 4 quartic terms of $V(\Sigma)$, and 3 independent mass terms from $V(\Sigma) + V(T, \Sigma)$ with T replaced by its VEV. The most general renormalizable potential for two Higgs doublets includes three more quartic terms. The allowed potential in our model is a special case of the general Two-Higgs-Doublet model, where $\lambda_1 = \lambda_2 = \lambda_3 \equiv \lambda_\Sigma$ and $\lambda_6 = \lambda_7 \equiv \tilde{\lambda}''_\Sigma$, using the standard notation [11]. The usual $-m_{12}^2 \tilde{H}_1 H_2$ mass mixing term arises here from the $\text{Tr}(\tilde{\Sigma}^\dagger \Sigma)$ term in $V(\Sigma)$ and the last term in $V(T, \Sigma)$, so that

$$m_{12}^2 = \text{Re} \left(\tilde{\lambda}'_{T\Sigma} u_T^2 - \tilde{M}_\Sigma^2 \right). \quad (2.13)$$

The squared masses of the charged Higgs boson and of the CP-odd scalar take a simple form:

$$\begin{aligned}
M_{H^\pm}^2 &= \frac{2m_{12}^2}{\sin 2\beta} - \left(\tilde{\lambda}_\Sigma + \tilde{\lambda}'_\Sigma \right) v_H^2 \quad , \\
M_A^2 &= M_{H^\pm}^2 + \left(\tilde{\lambda}_\Sigma - \tilde{\lambda}'_\Sigma \right) v_H^2 \quad .
\end{aligned}
\tag{2.14}$$

The value of m_{12} given by Eq. (2.13) may be comparable to $M_{W'}$, as both are controlled by the $SU(2)_R$ breaking VEV, u_T . At the same time, the weak scale is an order of magnitude smaller than u_T , which requires some tuning of the parameters in the potential; a similar tuning could lead to $m_{12} \ll M_{W'}$. Thus, theoretically M_{H^\pm} and M_A may be anywhere between the weak scale and $\sim M_{W'}$.

Various searches for the heavy Higgs bosons set mass limits substantially above the mass of the SM-like Higgs boson ($M_h = 125$ GeV). It is sufficient then to expand in $(v_H/m_{12})^2 \ll 1$. The heavy CP-even scalar, H^0 , has a squared mass

$$M_{H^0}^2 = M_{H^\pm}^2 + \left(\tilde{\lambda}_\Sigma + \tilde{\lambda}'_\Sigma \right) \cos^2 2\beta v_H^2 \quad .
\tag{2.15}$$

Note that the relative mass splittings between H^\pm , A^0 and H^0 are small. These four states approximately form a heavy $SU(2)_W$ doublet of zero VEV.

The mass of the SM-like Higgs boson is related to the quartic couplings by

$$M_h^2 = \left[\left(\tilde{\lambda}_\Sigma + \tilde{\lambda}'_\Sigma \right) \sin^2 2\beta + 2\lambda_\Sigma \right] v_H^2 \quad .
\tag{2.16}$$

This provides an estimate of the typical values of the quartic couplings: for $\tan \beta \rightarrow 1$, $2\lambda_\Sigma + \tilde{\lambda}_\Sigma + \tilde{\lambda}'_\Sigma \approx 0.5$. The departure from the alignment limit is given by

$$\alpha - \beta + \frac{\pi}{2} \simeq -\frac{1}{2} \left(\tilde{\lambda}_\Sigma + \tilde{\lambda}'_\Sigma \right) \sin 4\beta \frac{v_H^2}{M_{H^\pm}^2} \quad .
\tag{2.17}$$

3 W' decays into heavy Higgs bosons

Given the Higgs sector discussed in the previous section, which breaks the $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge symmetry down to the $U(1)_{\text{em}}$ group of electromagnetism, we can now derive the W' couplings to bosons and the ensuing decay widths.

3.1 Couplings of the W' to bosons

The $SU(2)_L$ and $SU(2)_R$ gauge bosons include electrically-charged states, $W_L^{\pm\mu}$ and $W_R^{\pm\mu}$, respectively. The kinetic terms for the T and Σ scalars give rise to the following mass terms for the charged gauge boson

$$(W_L^{+\mu}, W_R^{+\mu}) \begin{pmatrix} g_{2L}^2 \frac{v_H^2}{2} & -g_{2L}g_{2R} \frac{v_H^2}{2} \sin 2\beta \\ -g_{2L}g_{2R} \frac{v_H^2}{2} \sin 2\beta & g_{2R}^2 \left(u_T^2 + \frac{v_H^2}{2} \right) \end{pmatrix} \begin{pmatrix} W_{L\mu}^- \\ W_{R\mu}^- \end{pmatrix}, \quad (3.1)$$

where g_{2L} and g_{2R} are the $SU(2)_L \times SU(2)_R$ gauge couplings. Since we are interested in the case $M_{W'} \gg M_W$, we diagonalize the above mass matrix by keeping only the leading term in $(M_W/M_{W'})^2$. All equations that follow are valid up to corrections of order $(M_W/M_{W'})^2 \approx 0.2\%$, where we used $M_{W'} \approx 1.9$ TeV. The $SU(2)_L$ gauge coupling is given by the SM $SU(2)_W$ gauge coupling g , while the $SU(2)_R$ gauge coupling is given by the W' coupling g_R to the $u\bar{d}$ quarks in the gauge eigenstate basis:

$$\begin{aligned} g_{2L} &= g \approx 0.65, \\ g_{2R} &= g_R \approx 0.45-0.6, \end{aligned} \quad (3.2)$$

where the range for g_R is required in order to explain [6] the LHC mass peaks near 2 TeV.

The physical bosons, W and W' , are admixtures of $W_L^{\pm\mu}$ and $W_R^{\pm\mu}$:

$$\begin{aligned} W_\mu^\pm &= W_{L\mu}^\pm \cos \theta_+ + W_{R\mu}^\pm \sin \theta_+, \\ W'_\mu^\pm &= -W_{L\mu}^\pm \sin \theta_+ + W_{R\mu}^\pm \cos \theta_+. \end{aligned} \quad (3.3)$$

The $W_L^{\pm\mu} - W_R^{\pm\mu}$ mixing angle θ_+ satisfies

$$\sin \theta_+ = \frac{g_R}{g} \left(\frac{M_W}{M_{W'}} \right)^2 \sin 2\beta, \quad (3.4)$$

and the W and W' masses are given by

$$\begin{aligned} M_W &= \frac{g v_H}{\sqrt{2}}, \\ M_{W'} &= g_R u_T. \end{aligned} \quad (3.5)$$

For $M_{W'} \approx 1.9$ TeV and $g_R \approx 0.45-0.6$ (as determined in [6], by comparing the W' production cross section to the CMS dijet excess [4]), we find the $SU(2)_R$ breaking scale

$u_T \approx 3\text{--}4$ TeV. This set of parameters is compatible with constraints from electroweak precision observable [13].

The $W'WZ$ interactions are given by

$$\frac{g}{c_W} \sin \theta_+ i \left[W_\mu'^+ \left(W_\nu^- \partial^{[\nu} Z^{\mu]} + Z_\nu \partial^{[\mu} W^{-\nu]} \right) + Z_\nu W_\mu^- \partial^{[\nu} W'^{+\mu]} \right] + \text{H.c.} \quad (3.6)$$

Here $c_W \equiv \cos \theta_W$ is the usual SM parameter, and $[\mu, \nu]$ represents commutation ($\mu\nu - \nu\mu$) of the two Lorentz indices. Both $W_L^+ W_L^- Z$ and $W_R^+ W_R^- Z$ terms contribute to the above $W'WZ$ interactions through W_L - W_R mixing, with coupling strengths gc_W and gs_W^2/c_W , respectively. Note that the coefficient of Eq. (3.6) can also be written as $g_R (M_W/M_{W'})^2 \sin 2\beta/c_W$. Comparing this form with Eq. (4) of [6] gives $\xi_Z = \sin 2\beta$.

The W' interactions with a W and a neutral Higgs boson arise from the Σ kinetic term:

$$-g_R M_W W_\mu'^+ W^{\mu-} \left[h^0 \cos(\alpha + \beta) + H^0 \sin(\alpha + \beta) - iA^0 \cos 2\beta \right] + \text{H.c.} \quad (3.7)$$

also a W'^{\pm} interaction with ZH^\mp , which gets contributions from the kinetic terms of Σ and T :

$$-\frac{g_R}{c_W} M_W \cos 2\beta W_\mu'^{\pm} Z^\mu H^\mp . \quad (3.8)$$

The W'^{\pm} couplings to H^\mp and one of the neutral Higgs bosons are the following:

$$\frac{ig_R}{2} W_\mu'^+ \left[-\cos(\alpha + \beta) (H^- \partial^\mu H^0 - H^0 \partial^\mu H^-) + \sin(\alpha + \beta) (H^- \partial^\mu h^0 - h^0 \partial^\mu H^-) \right] + \text{H.c.} \quad (3.9)$$

In the alignment limit with $\tan \beta \rightarrow 1$ the above $W'^{\pm} H^\mp h^0$ coupling vanishes, while the $W'^{\pm} H^\mp H^0$ coupling reaches its maximum. Finally, the $W'^{\pm} H^\mp A^0$ coupling is

$$\frac{g_R}{2} \sin 2\beta W_\mu'^+ (H^- \partial^\mu A^0 - A^0 \partial^\mu H^-) + \text{H.c.} \quad (3.10)$$

3.2 W' branching fractions

The dominant decay modes of W' are into fermion pairs, mainly quark-antiquark pairs. W' can also decay into bosons, leading to challenging and interesting phenomenology. Amongst many possible bosonic decay modes, $W' \rightarrow WZ$ and $W' \rightarrow Wh^0$ are inevitable, originating from the kinetic term of the bidoublet scalar field Σ . Unlike other bosonic

decays discussed below, the widths of these two modes do not depend on the unknown masses of heavy scalars. Eqs. (3.6) and (3.7) imply

$$\begin{aligned}\Gamma(W' \rightarrow WZ) &\simeq \frac{g_{\text{R}}^2}{192\pi} \sin^2 2\beta M_{W'} \quad , \\ \Gamma(W' \rightarrow Wh^0) &\simeq \frac{g_{\text{R}}^2}{192\pi} \cos^2(\alpha + \beta) M_{W'} \quad ,\end{aligned}\tag{3.11}$$

where terms of order $(M_{W,Z,h}/M_{W'})^2$ have been neglected, and the approximate relation $c_W M_Z \simeq M_W$ has been used. We use full expressions for numerical study later in this section, which lead to a slight difference between the partial widths of these two W' decay modes, arising from differences in underlying dynamics and kinematics.

The equivalence theorem requires that W' decays into fields that are part of the same Higgs doublet (the longitudinal Z and h^0 for $\alpha \rightarrow \beta - \pi/2$ in this case) have equal decay widths up to electroweak symmetry breaking effects and phase-space factors. As mentioned in Ref. [6], in the alignment limit $\Gamma(W' \rightarrow WZ) \simeq \Gamma(W' \rightarrow Wh^0)$.

Besides h^0 and the longitudinal W and Z , the bidoublet field Σ includes the heavy scalars H^\pm, H^0, A^0 . The range of allowed masses for these particles spans more than an order of magnitude, from the weak scale to the $SU(2)_R$ breaking scale. If they are lighter than the W' boson, then W' decays may provide the main mechanisms for production of these scalar particles at hadron colliders. The W' decays into a heavy scalar and an electroweak boson have widths

$$\begin{aligned}\Gamma(W'^{\pm} \rightarrow ZH^\pm) &\simeq \frac{g_{\text{R}}^2}{192\pi} \cos^2 2\beta M_{W'} \left(1 - \frac{M_{H^\pm}^2}{M_{W'}^2}\right)^3 \quad , \\ \Gamma(W' \rightarrow W\phi) &\simeq \frac{g_{\text{R}}^2}{192\pi} \xi_\phi^2 M_{W'} \left(1 - \frac{M_\phi^2}{M_{W'}^2}\right)^3 \quad ,\end{aligned}\tag{3.12}$$

where ϕ labels the heavy neutral scalars, and ξ_ϕ follows from the couplings in Eq. (3.7):

$$\xi_\phi^2 = \begin{cases} \sin^2(\alpha + \beta) & , \quad \text{for } \phi = H^0 \quad , \\ \cos^2 2\beta & , \quad \text{for } \phi = A^0 \quad . \end{cases}\tag{3.13}$$

Here we neglected terms of order $M_W/M_{W'}$, which are relevant only for $M_{H^\pm} + M_W$ close to $M_{W'}$. The exact expressions are given by replacing the last factor in Eq. (3.12),

$$\left(1 - \frac{M_\phi^2}{M_{W'}^2}\right)^3 \rightarrow F\left(\frac{M_\phi^2}{M_{W'}^2}, \frac{M_W^2}{M_{W'}^2}\right) \quad ,\tag{3.14}$$

and a similar replacement for $W'^{\pm} \rightarrow ZH^{\pm}$, where the function F is defined by

$$F(x, y) = f^3(x, y) + 8y(1 + 2xy) f(x, y) \quad (3.15)$$

and f is the phase-space suppression factor:

$$f(x, y) = \left(1 - 2(x + y) + (x - y)^2\right)^{1/2} . \quad (3.16)$$

Similarly, the W' decay width into a charged Higgs boson and the SM-like Higgs particle is

$$\Gamma(W'^{\pm} \rightarrow H^{\pm}h^0) = \frac{g_R^2}{192\pi} \sin^2(\alpha + \beta) M_{W'} \left(1 - \frac{M_{H^{\pm}}^2}{M_{W'}^2}\right)^3 , \quad (3.17)$$

The terms of order $M_h/M_{W'}$, neglected here, are taken into account by replacing the last factor in the above equation by $f^3(M_{H^{\pm}}^2/M_{W'}^2, M_h^2/M_{W'}^2)$. For $M_{H^{\pm}} < M_{W'}/2$, the W' can also decay into a pair of heavy scalars, with widths

$$\Gamma(W'^{\pm} \rightarrow H^{\pm}\phi) = \frac{g_R^2}{192\pi} (1 - \xi_{\phi}^2) M_{W'} \left(1 - \frac{4M_{H^{\pm}}^2}{M_{W'}^2}\right)^{3/2} , \quad (3.18)$$

where ξ_{ϕ}^2 is defined in Eq. (3.13). Here we used the $M_A = M_{H^0} = M_{H^{\pm}}$ limit; the exact expression is obtained by replacing the last factor with $f^3(M_{H^{\pm}}^2/M_{W'}^2, M_{\phi}^2/M_{W'}^2)$.

The 8 widths for W' decays into bosons shown in this section satisfy the equivalence theorem for $M_{H^{\pm}}^2 \ll M_{W'}^2$ in the alignment limit. Summing over these 8 widths gives

$$\Gamma(W' \rightarrow \text{bosons}) = \frac{g_R^2}{48\pi} M_{W'} \left[1 - O(4M_{H^{\pm}}^2/M_{W'}^2)\right] . \quad (3.19)$$

Note that the leading order in $M_{H^{\pm}}^2/M_{W'}^2$ is independent of $\tan\beta$.

The dominant decay modes of W' are into quarks, and have the following widths:

$$\Gamma(W' \rightarrow jj) = \frac{g_R^2}{8\pi} M_{W'} \quad (3.20)$$

for light flavors, and

$$\Gamma(W'^+ \rightarrow t\bar{b}) = \frac{g_R^2}{16\pi} M_{W'} \left(1 + \frac{m_t^2}{2M_{W'}^2}\right) \left(1 - \frac{m_t^2}{M_{W'}^2}\right)^2 \quad (3.21)$$

for heavy flavors. QCD corrections increase these two widths by about 3%. The decay widths into a τ lepton or an electron and the N^{τ} right-handed neutrino are

$$\Gamma(W' \rightarrow \tau N^{\tau}) = \frac{g_R^2}{48\pi} (1 - s_{\theta_{\ell}}^2) M_{W'} \left(1 + \frac{m_{N^{\tau}}^2}{2M_{W'}^2}\right) \left(1 - \frac{m_{N^{\tau}}^2}{M_{W'}^2}\right)^2 ,$$

$$\Gamma(W' \rightarrow e N^{\tau}) = \frac{s_{\theta_{\ell}}^2}{1 - s_{\theta_{\ell}}^2} \Gamma(W' \rightarrow \tau N^{\tau}) , \quad (3.22)$$

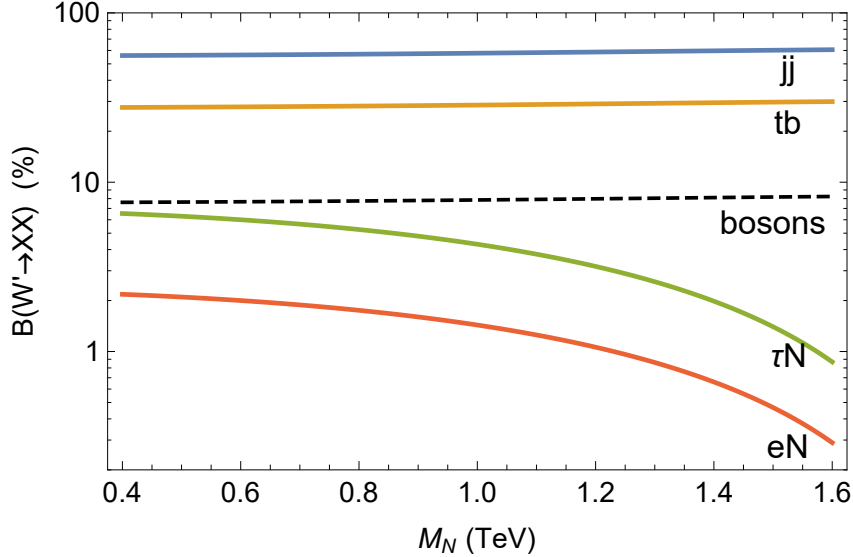


Figure 1: Branching fractions of W' for $M_{W'} = 1.9$ TeV and A^0 , H^0 , H^\pm masses of 500 GeV. The dashed line represents the sum of all 8 bosonic decay modes. The $W' \rightarrow \tau N^\tau$ and $e N^\tau$ widths are computed for $s_{\theta_e} = 0.5$.

where s_{θ_e} is the coefficient of the $g_R W'_\nu \bar{e}_R \gamma^\nu N_R^\tau$ interaction term in the Lagrangian, and the analogous coefficient for the muon satisfies $s_{\theta_\mu} \ll s_{\theta_e}$. The baseline W' model used in [6] to explain the excess events near 2 TeV reported in several channels by ATLAS and CMS, including e^+e^-jj [2], has $s_{\theta_e} \approx 0.5$, and is consistent with all current constraints on flavor-changing processes.² We emphasize that taking into account the e^+e^-jj excess was crucial in identifying our baseline W' model (without it, leptophobic W' models [15] are interesting alternatives). If the W' boson will be discovered in Run 2, then the $eejj$ process will allow various tests of the underlying couplings [16].

The $W' \rightarrow jj$ and $W' \rightarrow tb$ channels have a combined branching fraction of approximately 86%, as shown in Figure 1. The branching fractions for $W' \rightarrow \tau N_R^\tau$ and $W' \rightarrow e N_R^\tau$ add up to 6% for the benchmark value of $M_{N_R^\tau} = 1$ TeV. The remaining decay modes are into bosons, which together have a branching fraction of approximately 8% for $M_{H^\pm} = 500$ GeV, as indicated by the dashed line in Figure 1. Note that the difference between $\tan\beta = 1$ and 2 cannot be resolved, as expected from Eq. (3.19).

Eqs. (3.11)-(3.22) show that the W' widths into bosons depend on the almost degenerate masses of the heavy scalars, and on $\sin 2\beta$ whose value can be between 0.8 and 1,

²The tension with the $\mu \rightarrow e\gamma$ limit mentioned in [14] is not a concern for our model given that $s_{\theta_\mu} \rightarrow 0$ is a natural limit.

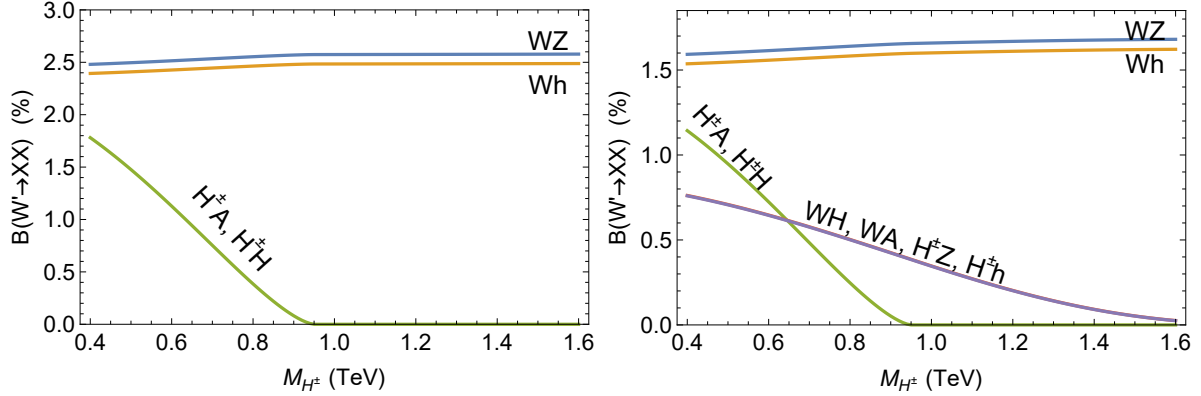


Figure 2: Branching fractions for W' decays into bosons, for $\tan\beta = 1$ (left panel) and $\tan\beta = 2$ (right panel). The mass of the N^τ right-handed neutrino is fixed at 1 TeV, and $M_{W'} = 1.9$ TeV. The $H^\pm A^0$ and $H^\pm H^0$ branching fractions are equal and not summed here (similarly for WA^0 , WH^0 , $H^\pm Z$ and $H^\pm h^0$).

as favored by the CMS and ATLAS mass peaks near $M_{W'} = 1.9$ TeV discussed in [6]. The dependence on the CP-even Higgs mixing angle α is very weak, due to the alignment limit Eq. (2.17), which implies $\cos(\alpha + \beta) \simeq \sin 2\beta$ and $\cos(\alpha - \beta) \ll 1$. We show the W' branching fractions in Figure 2 as a function of $M_{H^\pm} \simeq M_A \simeq M_{H^0}$, for $\sin 2\beta = 1$ (*i.e.*, $\tan\beta = 1$) and $\sin 2\beta = 0.8$ (*i.e.*, $\tan\beta = 2$ or $1/2$). The W' branching fractions to Wh^0 and WZ are almost constant because these partial widths do not depend on the masses of the heavy scalars. The branching fractions of W' into $H^\pm H^0$ and $H^\pm A^0$ are equal because of the alignment condition and the approximate mass degeneracy. Similarly, decay widths of WH^0 , WA^0 , $H^\pm Z$ and $H^\pm h^0$ are equal, suppressed by $\cos^2 2\beta$, and with a different phase-space function than the two-heavy-scalar modes. We can see from Figure 2 that the $\tan\beta = 1$ case is simple: WZ , Wh^0 , $H^\pm H^0$ and $H^\pm A^0$ modes are maximal, and the other decay modes vanish. When $\tan\beta$ deviates from unity, the other four decay modes start to emerge with sub-percent level branching fractions.

4 Quark masses

The $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge symmetry allows quark masses to be generated by Yukawa couplings to the bidoublet Σ :

$$-\bar{Q}_L^i \left(y_{ij} \Sigma + \tilde{y}_{ij} \tilde{\Sigma} \right) Q_R^j + \text{H.c.} \quad (4.1)$$

Here $i, j = 1, 2, 3$ label the SM fermion generations, $Q_L^i = (u_L^i, d_L^i)$ is the $SU(2)_L$ quark doublet of the i th generation, and $Q_R^i = (u_R^i, d_R^i)$ is the corresponding $SU(2)_R$ doublet. The Yukawa couplings y_{ij} and \tilde{y}_{ij} are complex numbers. The mass terms for the up- and down-type quarks are then

$$- v_H \bar{u}_L^i (y_{ij} \cos\beta + \tilde{y}_{ij} \sin\beta) u_R^j - v_H \bar{d}_L^i (y_{ij} \sin\beta + \tilde{y}_{ij} \cos\beta) d_R^j + \text{H.c.} \quad (4.2)$$

These terms highlight a problem. As we are interested in $\tan\beta = O(1)$, the above terms generically induce masses of the same order of magnitude for up- and down-type quarks. Even though the fermion mass hierarchies are not understood in the SM, they can be fitted by appropriately small Yukawa couplings. In the case of Eq. (4.2), the known ratio of the bottom and top quark masses can be fitted only by tuning various parameters. For example, $m_b \ll v_H$ can be achieved by imposing $\tilde{y}_{33}/y_{33} \simeq -\tan\beta$. In the absence of an explanation for the tuning of these independent parameters, it is useful to explore alternative mechanisms for fermion mass generations.

Consider the case where y_{ij} and \tilde{y}_{ij} are negligibly small. A different type of gauge-invariant operator that can generate the quark masses is

$$- \frac{C_{ij}}{u_T^2} \bar{Q}_L^i \tilde{\Sigma} T^\dagger T Q_R^j + \text{H.c.} \quad , \quad (4.3)$$

where the $SU(2)_R$ indices of the scalar triplet T are contracted with T^\dagger such that $T^\dagger T$ belongs to the triplet representation of $SU(2)_R$. The flavor-dependent coefficients C_{ij} are complex dimensionless parameters, and form a 3×3 mass matrix C . We assume that the analogous operator with $\tilde{\Sigma}$ replaced by Σ has negligibly small coefficients. Replacing T by its VEV given in Eq. (2.3), the dimension-6 operator of Eq. (4.3) generates a 3×3 mass matrix for the up-type quarks,

$$\mathcal{M}_u = v_H \sin\beta C \quad , \quad (4.4)$$

while the down-type quarks remain massless at this stage.

There is, however, an analogous effective operator with T replaced by $\tilde{T} \equiv \sigma_2 T^* \sigma_2$,

$$- \frac{\tilde{C}_{ij}}{u_T^2} \bar{Q}_L^i \tilde{\Sigma} \tilde{T}^\dagger \tilde{T} Q_R^j + \text{H.c.} \quad (4.5)$$

This generates a mass matrix only for down-type quarks:

$$\mathcal{M}_d = v_H \cos\beta \tilde{C} \quad . \quad (4.6)$$

Given that the coefficients \tilde{C}_{ij} are different than C_{ij} , we obtain mass matrices for the up- and down-type quarks that are independent of each other, as in the SM. Various possibilities for the origin of the effective operators (4.3) and (4.5) will be discussed elsewhere.

The operators (4.3) and (4.5) have an additional useful property: they separate the contributions of the two Higgs doublets to the quark masses. Notice that the embedding (2.12) of H_1 and H_2 in Σ implies $\tilde{\Sigma} = (\tilde{H}_2, H_1)$, where $\tilde{H}_2 = i\sigma_2 H_2^*$. From Eqs. (4.3) and (4.5) then follows

$$-C_{ij} \bar{Q}_L^i \tilde{H}_2 u_R^j - \tilde{C}_{ij} \bar{Q}_L^i H_1 d_R^j + \text{H.c.} \quad (4.7)$$

We have obtained the Yukawa couplings of the Two-Higgs-doublet model of Type II. This automatically avoids tree-level flavor-changing processes (see [17] and references therein). The couplings of the heavy Higgs bosons to the physical eigenstates of the quarks are proportional to the quark masses, with overall coefficients of $\sin\beta$ for up-type quarks and $\cos\beta$ for down-type quarks. For $0.5 \lesssim \tan\beta \lesssim 2$ and close to the alignment limit, as indicated [6] by the W' signals near 2 TeV, the branching fractions for $H^+ \rightarrow t\bar{b}$ and $H^0, A^0 \rightarrow t\bar{t}$ are almost 100%.

5 Signals of heavy scalars produced in W' decays

Even though the W' decays into heavy scalars have only percent-level branching fractions (see Figure 2), the large W' production cross section at the LHC [6] (200–350 fb at $\sqrt{s} = 8$ TeV, and 1–2 pb at $\sqrt{s} = 13$ TeV) makes these decays promising discovery channels. The $pp \rightarrow W' \rightarrow H^\pm A^0$ and $H^\pm H^0$ processes lead to cascade decays that end up with $3W + 4b$ final states, including

$$pp \rightarrow W'^+ \rightarrow H^+ A^0 / H^0 \rightarrow t\bar{b} t\bar{t} \rightarrow W^+ W^+ W^- + 4b \rightarrow \begin{cases} \ell^+ \ell^+ jj + 4b + \cancel{E}_T \\ \text{or} \\ \ell^+ \ell^+ \ell^- + 4b + \cancel{E}_T \end{cases} \quad (5.1)$$

where $\ell = e$ or μ . The charge conjugate processes lead to a number of events smaller by a factor of almost 2, due to the smaller W^- production in pp collisions. The cross section times branching fraction relevant for these processes is shown in Figure 3 for the LHC at 8 and 13 TeV. We used there $\tan\beta = 1$ and a W' production cross section of 300 fb at $\sqrt{s} = 8$ TeV for $M_{W'} = 1.9$ TeV (implying a cross section of 1.7 pb at $\sqrt{s} = 13$ TeV).

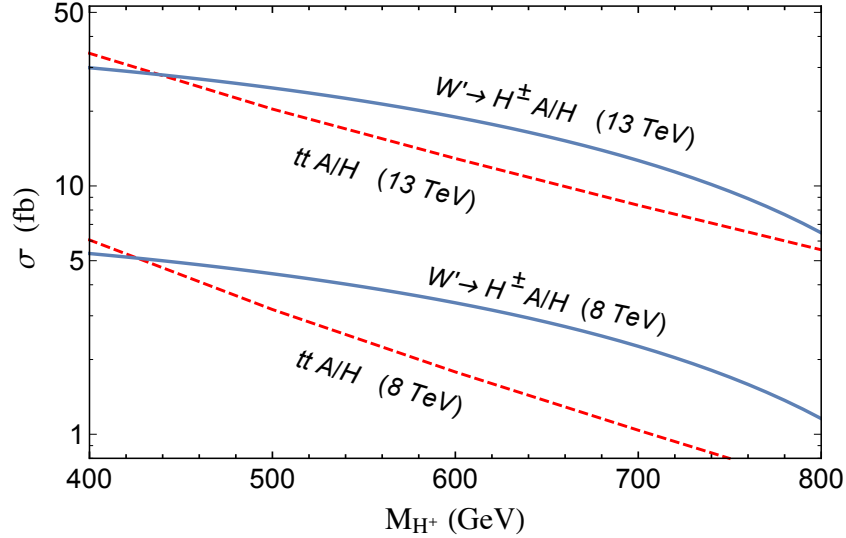


Figure 3: Production cross section times branching fractions for the sum of $pp \rightarrow W' \rightarrow H^\pm A^0, H^\pm H^0$ (solid lines), and for the sum of $pp \rightarrow t\bar{t}A^0, t\bar{t}H^0$ (dashed lines), at $\sqrt{s} = 8$ and 13 TeV. The parameters used here are $\tan \beta = 1$, $M_{W'} = 1.9$ TeV and a W' production cross section of 300 fb at 8 TeV.

An additional W' decay mode that contributes to final states with same-sign leptons or three leptons and b jets is

$$pp \rightarrow W'^+ \rightarrow \tau^+ N^\tau \rightarrow \tau^+ \tau^- t\bar{b} \rightarrow \ell^+ \nu \bar{\nu} \tau^- W^+ b\bar{b} \rightarrow \begin{cases} \ell^+ \ell^+ \tau_h^- b\bar{b} + \cancel{E}_T \\ \text{or} \\ \ell^+ \ell^+ \ell^- b\bar{b} + \cancel{E}_T \end{cases} \quad (5.2)$$

The branching fraction for the N^τ decay used here is $B(N^\tau \rightarrow \tau^- t\bar{b}) \approx 23\%$, while similar final states arise from $B(N^\tau \rightarrow e^- t\bar{b}) \approx 8\%$. The cascade decays $W'^+ \rightarrow e^+ N^\tau$ with $N^\tau \rightarrow \tau t\bar{b}$ or $e t\bar{b}$ (the latter process is mentioned in [18]) also contribute.

The heavy scalars may also be produced directly, without W' decays. A promising channel is the production of A^0 and H^0 in association with a $t\bar{t}$ pair [19]:

$$pp \rightarrow t\bar{t}A^0/H^0 \rightarrow 4t \rightarrow 4W + 4b \quad . \quad (5.3)$$

Compared to the SM $t\bar{t}h^0$ process at large M_h , each of the above two processes has a cross section scaled by $(\tan \beta)^{-2}$. In Figure 3 we show the cross section for this process, computed at leading order using MadGraph5_aMC@NLO [20] with the MSTWnlo2008 [21] parton distribution functions. We see that the processes listed in Eqs. (5.1) and (5.3) have comparable cross section times branching fractions.

The ATLAS search [22] for same-sign leptons and b -jets shows some deviations from the SM in two signal regions designed for four-top final states, with a pair of same-sign leptons and $H_T > 700$ GeV. In the signal region “SR4t3” (exactly two b -jets and $\cancel{E}_T > 100$ GeV) there are 12 events with an expected background of $4.3 \pm 1.1 \pm 1.1$ events. In the signal region “SR4t4” (3 or more b -jets) the search found 6 events with an expected background of $1.1 \pm 0.9 \pm 0.4$ events. We estimate that the combination of these signals represents a $\sim 3\sigma$ excess over the SM background.

The processes shown in Eqs. (5.1)-(5.3) provide a possible origin of this excess. The various contributions of heavy scalars and the right-handed neutrino to the two signal regions are shown in Table 1 for $\tan\beta = 1$ in the upper block and $\tan\beta = 2$ in the lower block. We set $M_{H^\pm} = M_{H^0} = M_{A^0} = 500$ GeV and 100% branching fractions of the heavy scalars into $t\bar{t}$ or tb . We estimate the efficiency by folding in the branching fractions to contributing final states, including combinatorial factors, and b -tagging efficiencies (using 70% for a single b tag). To obtain the predicted number of signal events in 20 fb^{-1} of data, we use a rough estimate of the acceptance, $A = 50\%$. For both values of $\tan\beta$, the expected number of signal events is compatible for both signal regions with the observed deviations above the background. We note that the W' -produced events include more leptons with positive charge than with negative charge; the same feature is seen in the ATLAS events (Tables 11 and 12 of [22]).

If the heavy scalar masses are decreased to $M_{H^\pm} \approx 400$ GeV, the number of signal events in Table 1 increases by a factor of about 3, so that the number of predicted events becomes too large (at $\tan\beta = 2$ it may still be acceptable, given the uncertainties in the event selection efficiencies). For $M_{H^\pm} \gtrsim 700$ GeV the number of predicted events becomes too small to account for the ATLAS excess. Thus, the preferred mass range for the heavy Higgs bosons is 400–700 GeV.

A CMS search [23] in a similar final state with same-sign leptons and b jets has yielded a smaller excess. The sum over the number of events with two or more b jets, large H_T and high lepton p_T from Table 3 of [23] gives 11 observed events for a background of roughly 6 ± 2 events. The compatibility of the ATLAS and CMS results in these channels needs further scrutiny. We note though that the CMS event selection includes a veto of a third lepton, while the events reported by ATLAS include events with three leptons.³

³In addition, the CMS jets are required to have $p_T > 40$ GeV, while for the ATLAS search a jet $p_T > 25$ GeV cut is imposed. Note that the energy released in the W' decay is shared between 10 or more objects, so that the stronger jet p_T cut imposed by CMS may reduce the number of observed events.

Signal	channel	efficiency	signal events	obs. (background)	
$bb\ell^\pm\ell^\pm$	$H^\pm H^0, H^\pm A^0 \rightarrow 3t + b$	2.5×10^{-4}	1.0-1.8	4.2-6.5	12 ($4.3 \pm 1.1 \pm 1.1$)
	$\tau N^\tau, e N^\tau \rightarrow (\tau/e)(\tau/e)tb$	5.2×10^{-4}	2.1-3.7		
	$t\bar{t}A^0, t\bar{t}H^0 \rightarrow 4t$	1.6×10^{-2}	1.1		
$\geq 3b\ell^\pm\ell^\pm$	$H^\pm H^0, H^\pm A^0 \rightarrow 3t + b$	6.2×10^{-4}	2.5-4.4	5.1-7.0	6 ($1.1 \pm 0.9 \pm 0.4$)
	$t\bar{t}A^0, t\bar{t}H^0 \rightarrow 4t$	4.1×10^{-2}	2.6		
$bb\ell^\pm\ell^\pm$	$H^\pm H^0, H^\pm A^0 \rightarrow 3t + b$	1.6×10^{-4}	0.7-1.1	4.2-7.2	12 ($4.3 \pm 1.1 \pm 1.1$)
	$WH^0, WA^0 \rightarrow Wt\bar{t}$	2.2×10^{-4}	0.9-1.6		
	$H^\pm h, H^\pm Z \rightarrow \ell tb + X$	0.7×10^{-4}	0.3-0.5		
	$\tau N^\tau, e N^\tau \rightarrow (\tau/e)(\tau/e)tb$	5.2×10^{-4}	2.1-3.7		
	$t\bar{t}A^0, t\bar{t}H^0 \rightarrow 4t$	1.6×10^{-2}	0.3		
$\geq 3b\ell^\pm\ell^\pm$	$H^\pm H^0, H^\pm A^0 \rightarrow 3t + b$	4.0×10^{-4}	1.6-2.8	2.3-3.5	6 ($1.1 \pm 0.9 \pm 0.4$)
	$t\bar{t}A^0, t\bar{t}H^0 \rightarrow 4t$	4.1×10^{-2}	0.7		

Table 1: Contributions from W' cascade decays and $t\bar{t}A^0/H^0$ production to the same-sign leptons plus b -jets signals at the 8 TeV LHC. The last column gives the observed and expected number of events in the ATLAS search [22]. For the upper and lower blocks of the table $\tan\beta = 1$ and 2, respectively. The range of predicted signal events corresponds to the 200–350 fb range for the W' production rate. The parameters are fixed as follows: heavy scalar masses are 500 GeV, the N^τ mass is 1 TeV, $s_{\theta_e} = 0.5$, and $M_{W'} = 1.9$ TeV.

Searches for final states with 3 leptons may also test the presence of the heavy Higgs particles, as follows from the processes in Eqs. (5.1)-(5.3). The CMS search [24] for events with 3 leptons and one or more b jets has yielded a deficit of events compared to the SM prediction. In particular, the large H_T category with no hadronic τ decays, $\cancel{E}_T > 100$ GeV and no e^+e^- or $\mu^+\mu^-$ pairs includes a single observed event for an expected background of 5.5 ± 2.2 events (Table 3 of [24]). Let us estimate the contribution to this event category from the $3t + b$ final state produced in the W' cascade decays of Eq. (5.1). Of the 1–1.8 signal events from the first row of Table 1, only a fraction of $B(W \rightarrow \ell + \cancel{E}_T)/4 \approx 6.4\%$ would pass the criterion of 3 leptons and no e^+e^- or $\mu^+\mu^-$ pairs. Similar suppressions apply to the number of 3ℓ events contributed by the processes listed in the other rows

of Table 1. Thus, searches in 3-lepton final states do not appear to be able for now to differentiate between the SM and our theory with a W' and heavy Higgs bosons.

The event selection employed in the four-top ATLAS search [22] has a few additional categories, including those with lower H_T , one b -jet, lower \cancel{E}_T , which have not lead to significant excesses as the categories “SR4t3” and “SR4t4” discussed above. The heavy Higgs contributions from W' decays to these other categories are much smaller. For instance, the heavy W' renders the H_T large (> 700 GeV), with harder \cancel{E}_T (> 100 GeV). Due to the use of “loose” b -tagging in the search and the presence of four b quarks in the heavy Higgs signal, the number of events in the one b -jet categories is a factor of 3–5 smaller than in the two b -jet (“SR4t3”) and three-or-more b -jet (“SR4t4”) categories.

6 Conclusions and outlook

A W' boson of mass near 1.9 TeV with properties detailed in Ref. [6] provides a compelling explanation for the excess events reported by the ATLAS and CMS Collaborations [1]-[5] in the WZ , Wh^0 , e^+e^-jj , and jj channels. In this article we have shown that the gauge structure associated with that W' boson predicts the existence of heavy Higgs bosons of masses between the weak scale and a few TeV. We have derived the branching fractions for all W' decays, including six channels with heavy Higgs bosons (see Figures 1 and 2).

The main decay modes for the heavy Higgs bosons are $A^0, H^0 \rightarrow t\bar{t}$ and $H^\pm \rightarrow t\bar{b}$. If their masses are below $M_{W'}/2$, then the cascade decay $W' \rightarrow H^\pm A^0/H^0 \rightarrow 3t + b \rightarrow 3W + 4b$ has a branching fraction of up to 3% and provides a promising way for discovering all these particles. An excess of events, with a statistical significance of about 3σ , has been reported by the ATLAS Collaboration [22] in the final state with two leptons of same charge and two or more b jets. We have shown that this can be explained by the cascade decays of W' if the heavy Higgs bosons have masses in the 400–700 GeV range.

The $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ gauge theory presented here depends on only a few parameters, whose ranges are already determined by accounting for the deviations from the SM mentioned above. The various phenomena predicted by this gauge theory can thus be confirmed or ruled out in the near future. In Run 2 of the LHC, the W' production cross section is large, in the 1–2 pb range at $\sqrt{s} = 13$ TeV. Besides resonant production of WZ , Wh^0 , jj and tb , there are several W' discovery modes: $W' \rightarrow \tau N^\tau \rightarrow \tau\tau jj, \tau\tau tb, e\tau jj, e\tau tb$ and $W' \rightarrow e N^\tau \rightarrow ee jj, eetb$ would test the existence of the heavy right-handed neutrino,

while $W' \rightarrow H^\pm A^0/H^0 \rightarrow 3t + b$, $W' \rightarrow WA^0/H^0 \rightarrow Wt\bar{t}$, $W' \rightarrow H^\pm h^0 \rightarrow tbh^0$ and others would test the existence of the heavy Higgs bosons.

Another promising search channel for the heavy neutral Higgs bosons, independent of the W' , follows from production in association with a $t\bar{t}$ pair, which has a cross section of the order of 10 fb at $\sqrt{s} = 13$ TeV. With more data, the Z' boson analyzed in [6] will also be accessible in a variety of channels.

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