Associated production of the heavy charged gauge boson $W_H$ and a top quark at LHC

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Abstract – In the context of the topflavor seesaw model, we study the production of the heavy charged gauge boson $W_H$ associated with a top quark at the LHC. Focusing on the searching channel $pp \to tW_H \to t\bar{t}bH\nu\bar{\nu}jjb\bar{b}$, we carry out a full simulation of the signal and the relevant standard model backgrounds. The kinematical distributions of final states are presented. It is found that the backgrounds can be significantly suppressed by sets of kinematic cuts, and the signal of the heavy charged boson might be detected at the LHC with $\sqrt{s} = 14$ TeV. With an integrated luminosity of $L = 100$ fb$^{-1}$, a $3.8\sigma$ signal significance can be achieved for $m_{W_H} = 1.60$ TeV.

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Introduction. – The standard model (SM) of particle physics is one of the most successful theories over the past decades which describes a variety of experimental results. However, the theoretical shortcomings of the SM suggest that it should be embedded in a larger scheme. Many popular new physics models beyond the SM have been proposed, and some of which predict the existence of the new gauge bosons with masses at the TeV order. In this paper, we study the signal of the heavy charged gauge boson $W_H$ in the topflavor seesaw model at the LHC.

It is interesting to note that only the top quark mass lies at the same mass scale of weak gauge bosons, while all other fermions are provided with masses less than a few GeV. This suggests that the top quark sector may involve some new gauge dynamics at the weak scale in contrast to all light fermions. In the topflavor seesaw model, in order to address the mass hierarchy in fermions, only the top sector enjoys the extra $SU(2)_t$ gauge interaction which is stronger than the ordinary $SU(2)_f$ (associated with all the other light fermions). After the gauge symmetry breaking $SU(2)_t \times SU(2)_f \to SU(2)_L$, at a high scale $\Lambda$, the heavy gauge bosons which are combinations of the corresponding broken gauge fields of $SU(2)_t$ and $SU(2)_f$ obtain mass. This results in the fact that the coupling of $W_H$ with $t\bar{t}$ is enhanced by the mixing parameter $1/x$ while the couplings of $W_H$ with light fermions are suppressed by a factor of $x$ [1]. Here, $x$ is the ratio of the gauge coupling $g_1$ of $SU(2)_f$ to the gauge coupling $g_0$ of $SU(2)_t$ and it is constrained to be $x^2 \ll 1$ in this construction.

The feature of the topflavor seesaw model is different from many new physics models with heavy $W_H$ including extra dimension models [2], little Higgs theories [3] and several other well-motivated extensions of the SM [4–7]. In this paper, we study the production of $W_H$ associated with a top quark at the LHC. This process provides an independent test of the $W_Ht\bar{t}$ coupling which is closely relevant to the model feature.

This paper is organized as follows. In the second section, we give a brief review of the topflavor seesaw model. In the third section, the phenomenological analysis and numerical calculations for the production of the heavy charged gauge boson $W_H$ in association with a top quark are presented. Our conclusions are given in the last section.

A brief review of the topflavor seesaw model. – The topflavor seesaw model [8], in which the top sector experiences a new $SU(2)_t$ gauge interaction, is based on the gauge group $SU(3)_c \otimes SU(2)_L \otimes SU(2)_f \otimes U(1)_y$. The corresponding three gauge couplings of the gauge group $SU(2)_t \otimes SU(2)_f \otimes U(1)_y$ are $g_0, g_1, g_2$. The Higgs doublet $\Phi_1$ with a nonzero vacuum expectation value (VEV) $v$ breaks $SU(2)_t \otimes SU(2)_f$ down to the SM gauge group $SU(2)_L$, and the Higgs doublet $\Phi_2$ with a VEV $v$ breaks

\[SU(2)_t \otimes SU(2)_f \otimes U(1)_y \to SU(2)_L \otimes U(1)_y.\]

\[SU(2)_t \otimes SU(2)_f \otimes U(1)_y \to SU(2)_L \otimes U(1)_y.\]

\[SU(2)_t \otimes SU(2)_f \otimes U(1)_y \to SU(2)_L \otimes U(1)_y.\]
SU(2)_L \otimes U(1)_Y down to the SM gauge group U(1)_{em}. This kind of breaking pattern results in extra massive color-singlet heavy gauge bosons W_H and Z_H, in addition to the two neutral physical Higgs boson h^0 and H^0. It is exciting that the LHC has discovered a Higgs-like boson with mass around 125 GeV [9], which coincides with the light Higgs h^0 predicted by the topflavor seesaw model [1].

The structure of topflavor seesaw model including its Higgs, gauge and top sectors has been systematically studied in refs. [1,8]. In this work, we aim to study the production of the heavy charged boson W_H at the LHC. The couplings of the heavy charged boson W_H to ordinary particles, which are related to our calculation, are given by [1]

$$\mathcal{L}_{W_H^{\mp}f} = \frac{-ig}{\sqrt{2}} \gamma_{\mu} P_L V_{ij} W_H^{\mp} f_{\mu}^{ij} W_H^{\mp} + \text{h.c.},$$

(1)

and

$$\mathcal{L}_{H} = -ix(c_\alpha - y s_\alpha) g_{W_H} W^\mp_{\mu} W_{\mu} H^0 + ix(s_\alpha + y c_\alpha) g_{W_H} W_{\mu}^{\pm} W_{\mu}^{\mp} H^0,$$

where we have defined the ratios $x = g_t/g_b$, $r = k^2/M_Z^2$, and $y = v/F$. Here $M_S$ denotes the mass eigenvalues parameter of heavy quarks ($T, B$), and the parameter $k$ is expected to be of $O(M_S)$. And $\alpha$ represents the mixing angle between physical Higgs bosons ($h^0, H^0$) and their weak eigenstates. The Higgs-like particle found by the LHC agrees well with the SM prediction which suggests a small $\alpha$ in the range $0 < \alpha \leq 0.2\pi$ [1,9].

The coupling of $W_H$ with the light gauge bosons W and Z is given by

$$g_{W_H^WZ} = \frac{-ix^2 y^2}{c_W} G_{W_H^WZ}^{SM},$$

(3)

where $G_{W_H^WZ}^{SM}$ represents the coupling constant of WWZ in the SM.

A detailed analysis of direct search for the topflavor seesaw model at the LHC and the constraints on this model of electroweak precision measurements are presented in ref. [1]. In the topflavor seesaw model, the heavy spectator quarks ($T, B$) are vector-like under SU(2)_L, so the fermionic contributions to oblique corrections can be fairly small in the decoupling limit. The Higgs sector contributions are relevant to the masses of the physical Higgs bosons ($h^0, H^0$) and the mixing angle $\alpha$. The contributions from gauge sectors are dependent on the masses of the heavy gauge bosons $W_H/Z_H$ and the mixing parameter $x$. After deducing the contributions to electroweak precision parameters from gauge, Higgs and fermion sectors, it is found that at 95% CL, there should be $m_{W_H} \geq 0.45–1$ TeV for a wide $H^0$ mass range up to 800 GeV with the inputs $\alpha = 0.1\pi–0.2\pi$ and $M_T = 4$ TeV [1].

The ATLAS and CMS collaborations have been actively searching for the new gauge boson $W_H$ at the LHC [10,11]. They mainly focus on the sequential standard model (SSM), where the couplings of $W_H$ with fermions equal the corresponding SM couplings. The most stringent limit on the $W_H$ mass comes from $pp \rightarrow W_H \rightarrow tb$ given by the CMS collaboration [10]. In particular, the CMS detector illustrates the mass of right-handed $W_H$ should be larger than 0.84 TeV through the $t\bar{t}W_H$ associated production [12]. However, different scenarios have different phenomenological features. The direct search constraints can be relaxed in some models. In the topflavor seesaw model, the couplings of heavy bosons $W_H/Z_H$ with light fermions are suppressed by the small mixing angles, which are of the order of $O(x)$. However, the coupling of $W_Htb$ is enhanced by a factor $1/x$ in the topflavor seesaw model.

With the sample input $x = 0.15$, it is found a 95% CL lower limit on the $W_H$ mass, $m_{W_H} > 1.0$ TeV, from the CMS data [1,10].

**Numerical results and discussion.** Before studying the process $pp \rightarrow t\bar{t}W_H$, we firstly consider the possible decay modes of the heavy charged gauge boson $W_H$. Once produced at the LHC, the heavy boson $W_H$ can decay into $f\bar{f}$, $tb$, $W^0 h^0$, $W^0 H^0$ and $WZ$. The partial widths of $W_H$ decaying to a pair of fermions are

$$\Gamma(W_H \rightarrow f\bar{f}) = \frac{g^2 x^2}{16\pi} |V_{ff'}|^2 m_{W_H},$$

(4)

$$\Gamma(W_H \rightarrow tb) = \frac{g^2}{16\pi} \left[\frac{1 - r x^2}{x(1 + r)}\right]^2 m_{W_H} \times \left(1 - \frac{m_t^2}{m_{W_H}^2} - \frac{m_h^2}{2m_{W_H}^2} - \frac{m_h^4}{4m_{W_H}^4}\right).$$

(5)

Likewise, the partial widths of $W_H$ decaying to a $W$ boson and a Higgs are

$$\Gamma(W_H \rightarrow W h^0) = \frac{g^2 x^2 (c_\alpha - y s_\alpha)^2}{192\pi m_{W_H}^2} \times \left((m_W^2 + m_h^2 - m_{h^0}^2)^2 - 24 m_W^2 m_h^2 m_{h^0}^2\right),$$

(6)

$$\Gamma(W_H \rightarrow W H^0) = \frac{g^2 x^2 (s_\alpha + y c_\alpha)^2}{192\pi m_{W_H}^2} \times \left((m_W^2 + m_h^2 - m_{h^0}^2)^2 - 24 m_W^2 m_h^2 m_{h^0}^2\right).$$

(7)

Noting the coupling $g_{W_H^WZ} \propto (x^2 y^2)$, and $x^2, y^2 \ll 1$, we neglect the $W_H \rightarrow Z\gamma$ decay channel here. The total decay width is about 57.9 GeV (71.6 GeV) for a 1.3 TeV (1.6 TeV) $W_H$ in the case $(x, r, \alpha) = (0.20, 1, 0.13\pi)$. We present branching ratios (BRs) of $W_H$ with various mass values in fig. 1(a). It is found that the BR of $W_H \rightarrow tb$ is dominant and the BRs of other decay modes $W_H \rightarrow f\bar{f}$, $W_H \rightarrow W h^0$ and $W_H \rightarrow W H^0$ are tiny. This is because the coupling of $W_Htb$ is roughly proportional to $1/x$ while the other decay modes are suppressed by the relevant couplings. Furthermore, we show...
the cross-section for process \( pp \rightarrow tW_H \) (include both \( gb \rightarrow tW_H^+ \) and \( gb \rightarrow tW_H^- \)) as a function of the heavy boson mass \( m_{W_H} \) with \( \sqrt{s} = 14 \text{ TeV} \) in fig. 1(b). As shown in fig. 1(b), the cross-sections decrease as \( m_{W_H} \) increases, which are in the ranges of hundreds fb to several fb, for \( m_{W_H} = 1000 \text{ GeV} - 2000 \text{ GeV} \) and \( x = 0.15 - 0.25 \). The cross-section for process \( pp \rightarrow tW_H \) in the topflavor seesaw model is larger than that in SSM, top-philic W' model [6] and little Higgs models [13].

In this paper, we will focus on the \( tW_H \) production followed by the dominant decay channel \( W_H \rightarrow tb \). For the \( ttb \) signal, there are three search modes: the full hadronic mode, the dileptonic mode and the semileptonic mode. In this paper, we will focus on the semileptonic mode. There are some advantages in this channel: 1) The charged lepton can provide an efficient trigger and suppress the large QCD backgrounds. 2) There is no reconstruction difficulty of momenta of invisible neutrino as encountered in the dileptonic mode. 3) There is no large combinatorial problem of multi-jets final states as that in the full hadronic mode. 4) It has a large branching ratio. Here, we demand the hadronic decay of the antitop \( \bar{t} \rightarrow bW^- \rightarrow bjj \) and the leptonic decay of the top quark \( t \rightarrow bW^+ \rightarrow b l^+ \nu_l \). The corresponding signal processes are

\[
\begin{align*}
pp &\rightarrow gb \rightarrow tW_H^+ \rightarrow t\bar{b} \rightarrow W^- b\bar{b} \rightarrow j j b l^+ \nu bb, \quad (8) \\
pp &\rightarrow gb \rightarrow tW_H^- \rightarrow t\bar{b} \rightarrow W^+ b\bar{b} \rightarrow t l^+ \nu b j j b b. \quad (9)
\end{align*}
\]

Both electrons and muons are considered for the positive charged lepton in our analysis. The Feynman diagrams for the signal processes are shown in fig. 2.

In this letter, we consider two typical points with \( m_{W_H} = 1.3 \text{ TeV}, 1.6 \text{ TeV} \) and employ sets of kinematical cuts to reconstruct the signal. For such heavy \( W_H \), the top quark from its decay often has a momentum much larger than \( m_t \). In this case, the boosted top tagging technique could help to identify the highly boosted top and suppress the backgrounds [14,15]. Roughly speaking, the top jet radius is \( \Delta R_{top} \sim \frac{2m_t}{p_T} \) and the typical top taggers efficiency can reach about 40% (dependent on the transverse momentum of top and algorithm parameters) with a corresponding QCD jet mis-tag rate at percent level [15]. The phenomenological study of charged Higgs search for process \( pp \rightarrow tH^- \rightarrow t\bar{b} \) in full hadronic mode has shown that a similar signal can be extracted from large SM QCD backgrounds and with the help of boosted top tagger [16]. However, we only reconstruct the signal in the traditional reconstruction method in this letter and leave the detailed analysis with top tagging in later study.

For the \( bbbb j l l \nu \) signal, the SM backgrounds mainly come from the irreducible \( t\bar{t}b \) background and the reducible \( t\bar{t}j \) background, where the light jet \( j \) means the light-flavor quarks or gluons:

\[
\begin{align*}
pp &\rightarrow t\bar{t}b \rightarrow bbbjjl^+\nu, \quad (10) \\
pp &\rightarrow t\bar{t}j \rightarrow bbbjjl^+\nu. \quad (11)
\end{align*}
\]

The other SM background processes, e.g., \( Wjjjjj, WWjjj \) and \( WZjjj \), etc. can be dramatically suppressed by the kinematical cuts and therefore we neglected them here.

Fig. 1: (Color online) (a) The branching ratio of \( W_H \) as functions of the heavy boson mass \( m_{W_H} \) in the case \( (x, r, \alpha) = (0.20, 1, 0.13\pi) \); (b) the cross-section for the process \( pp \rightarrow tW_H \) as a function of the heavy boson mass \( m_{W_H} \) for \( r = 1 \) and three typical values \( x \) at the LHC with \( \sqrt{s} = 14 \text{ TeV} \).

Fig. 2: (Color online) The partonic level process for a \( W_H \) in association with a top quark production and decay in hadronic collisions.

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In our study, MadGraph/MadEvent [17] is employed to generate both the exclusive signal and background events at the parton level. And then we simulate detector resolution effects by smearing the lepton and jet energies according to the assumed Gaussian resolution parametrization
\[
\delta(E) = \frac{a}{E} \oplus b.
\] (12)
We take \(a = 5\%\) and \(b = 0.55\%\) for leptons, and take \(a = 100\%\) and \(b = 5\%\) for jets [18,19].

The basic acceptance cuts, referred to as cut I, are applied for the signal and background events,
\[
p_{Tj} \geq 25 \text{ GeV}, \quad p_{Tl} \geq 25 \text{ GeV}, \quad E_T \geq 25 \text{ GeV}, \quad |\eta_l| \leq 2.5, \quad |\eta_j| \leq 2.5, \quad \sqrt{(\phi_i - \phi_j)^2 + (\eta_i - \eta_j)^2} \geq 0.4.
\] (13)
Here, \(\eta_i(\phi_i)\) denotes the rapidity (azimuthal angle) of the related lepton (jet). \(p_{Tj}(p_{Tl})\) is the jet (lepton) transverse momentum, \(E_T\) is the missing transverse momentum from the invisible neutrino in the final state.

After the basic cuts to simulate the detector acceptance, we further employ sets of kinematical cuts to reduce the backgrounds based on the kinematical differences between the signal and backgrounds. The signal events consist of five jets in the final state. The jets from the heavy boson \(W_H\) decay tend to have larger \(p_T\) than the jets in the background events. We order the jets by their \(p_T\) and present the normalized \(p_T\) distributions for the leading jet and the second jet in the signal events and background events in fig. 3. The model parameters are set as \(m_{W_H} = 1.30\ \text{TeV}\) and \((x, r, \alpha) = (0.20, 1, 0.13\pi)\) in the signal process.

From fig. 3(a), we can see that the leading jet \((j_1)\) in the signal events has much harder \(p_T\) distributions than the jet in the SM background events. This is because the hardest jet in the signal is mainly the \(b\) jet or the daughter-jet of the highly boosted top from the \(W_H \rightarrow \bar{t}b\) decay. However, the top quarks in the SM backgrounds are mainly produced in the threshold region. Thus the jets from top quark decay in the backgrounds tend to be soft. This difference enables us to impose a hard cut on \(p_T\) to further suppress the SM backgrounds.
Similar to the leading jet, the second jet \((j_2)\) in the signal is harder than that in the backgrounds as displayed in fig. 3(b). Furthermore, the normalized \(H_t\) distributions, i.e., the scalar sum of the \(p_T\)’s for all the visible particles in the final state, are shown in fig. 4(a). The model parameters are set as \(m_{W_H} = 1.30\) TeV and \((x, r, \alpha) = (0.20, 1, 0.13)\pi\) in the signal process.

From fig. 3 and fig. 4(a), we can see that there are large kinematic differences between the signal and backgrounds. To purify the signal, a set of hard \(p_T\) cuts are further adopted for \(m_{W_H} = 1.30\) TeV, based on above analysis, as follows:

\[
p_T(j_1) \geq 250\text{ GeV}, \quad p_T(j_2) \geq 140\text{ GeV},
\]

\[
H_t = \sum_i p_T(j_i) + p_T(l^+^) \geq 800\text{ GeV}.
\]

(14a)

The \(p_T\) spectrum of the background is independent of \(m_{W_H}\) whereas the \(p_T\) spectrum of the signal is sensitive to \(m_{W_H}\). For \(m_{W_H} = 1.60\) TeV, a similar set of cuts

\[
p_T(j_1) \geq 350\text{ GeV}, \quad p_T(j_2) \geq 180\text{ GeV},
\]

\[
H_t = \sum_i p_T(j_i) + p_T(l^+^) \geq 1200\text{ GeV}
\]

(14b)

is adopted. These cuts in eq. (14) are referred to as cut II, and the efficiencies of these cuts are shown in table 1.

In order to suppress the \(t\bar{t}j\) background, it is crucial to identify the extra jet (denoted as \(j_{\text{extra}}\)) produced in association with the \(t\bar{t}\) pair as a \(b\) jet. Following the reconstruction scheme [6], we apply the minimal \(\chi^2\) template method which is based on the \(W\) boson and top quark masses to pick out the extra jet,

\[
\chi^2 = \frac{(m_t - m_{tb})^2}{\sigma_t^2} + \frac{(m_t - m_{jjj})^2}{\sigma_t^2} + \frac{(m_W - m_{jjj})^2}{\sigma_W^2}.
\]

(15)

The transverse momentum of the neutrino can be determined from the momentum imbalance in the transverse plane while the longitudinal momentum of the neutrino \((p_{\nu,L})\) can be derived from the \(W\)-boson on-mass-shell condition, \(m_W^2 = m_t^2\). It yields a twofold solution as

\[
p_{\nu,L} = \frac{1}{2p_{T}^{-T}} \left[ \left( m_W^2 + 2p_{T}^{-T} \cdot \vec{E}_{T}^{-} \right) p_{T}^{-L} \pm E_{T}^{-} \sqrt{ \left( m_W^2 + 2p_{T}^{-T} \cdot \vec{E}_{T}^{-} \right)^2 - 4p_{T}^{-T} \vec{E}_{T}^{-}^2} \right].
\]

(16)

when \((m_W^2 + 2p_{T}^{-T} \cdot \vec{E}_{T}^{-})^2 - 4p_{T}^{-T} \vec{E}_{T}^{-} \geq 0\). The value of \(p_{\nu,L}\) that best yields the known top mass is selected via \(m_{tb}^2 = m_t^2\). The ambiguity of the twofold solution is removed also by the minimal \(\chi^2\) requirement.

In fig. 4(b), we present the normalized \(p_T\) distributions of the extra jet for \(W_H = 1.30\) TeV and \((x, r, \alpha) = (0.20, 1, 0.13)\pi\). The extra jet in association with \(t\bar{t}\) in the SM backgrounds comes mainly from QCD radiation, while the extra jet in the signal is often predominately from the heavy \(W_H\) decay. Hence, it is obvious that the extra jet in the signal tends to have a much harder \(p_T\) distributions than the extra jet in the SM backgrounds. Here, we further take a cut on the extra jet,

\[
p_T(j_{\text{extra}}) \geq 200\text{ GeV}.
\]

(17)

For \(m_{W_H} = 1.60\) TeV, the larger mass is transferred to the momentum of final states. As stated above, the reconstructed extra jet is always the \(b\)-jet from heavy charged boson decay in signal events and the \(p_T\) of it is sensitive to the mass of the parent \(W_H\). So, \(p_T(j_{\text{extra}}) \geq 400\) GeV is adopted in the case \(m_{W_H} = 1.60\) TeV.

After the extra jet is discriminated, we further require it to be a \(b\)-jet to suppress the \(t\bar{t}j\) background. In order to keep more signal events, we only take this one \(b\)-tagging while not require tagging the \(b\)-jet from top decay. Here, the \(b\)-tagging efficiency of 60% is used and the mis-tagging efficiency of a light quark jet and gluon jet is taken as 1% [19].

After reconstructing both the top and antitop and singling out the extra jet, it is easy to reconstruct the \(W_H\). Since our signal events consist of not only \(t\bar{W}_H^-\) but also \(W_H^+\), both \(j_{\text{extra}}\) and \(l_{\text{extra}}\) might be the corrected reconstructed \(W_H\). So, we recombine the harder reconstructed top and extra jet to get the reconstructed \(W_H\) because the top decaying from the heavy charged gauge boson with TeV mass always carries large momentum. And we further require the invariant mass of the extra jet and reconstructed \(t\) quark or \(\bar{t}\) quark to be around the heavy boson \(W_H\) mass window,

\[
|m_{W_H} - m_{j_{\text{extra}}} | < 200\text{ GeV}, \quad \text{or} \quad |m_{W_H} - m_{l_{\text{extra}}} | < 200\text{ GeV}.
\]

(18)

The cuts in eq. (18) can efficiently suppress the SM backgrounds while keeping most of the signal. These cuts in eq. (17) and eq. (18) are referred to as cut III.

As shown in table 1, the sets of cuts significantly suppress the backgrounds. Supposing the integrated luminosity to be \(100\text{ fb}^{-1}\) at \(\sqrt{s} = 14\) TeV, a large significance \(S/\sqrt{B} = 13.9\) (8.32) can be achieved for \(W_H = 1.30\) TeV \((W_H = 1.60\) TeV\) with \((x, r, \alpha) = (0.20, 1, 0.13)\pi\).
Conclusions. — In this paper, focusing on the semileptonic channel for process $pp \rightarrow tW_H \rightarrow t\bar{b}$, we have studied the potential of discovering the extra heavy gauge boson $W_H$ predicted in the topflavor seesaw model at the LHC. Studying this process can independently test the $W_Htb$ coupling and shed light on the flavor structure and the gauge structure of new physics models. In the topflavor seesaw model, the couplings of $W_H$ with $\bar{b}t$ is enhanced by a factor of $1/x$ while the couplings of $W_H$ with light fermions are suppressed by a factor of $x$. After calculation, it is found that the cross-section for $tW_H$ production can reach tens of fb for $m_{W_H}$ in the mass range 1.2–1.7 TeV. We further studied the kinematic features of the signal and employ sets of cuts to extract the signal from backgrounds. Our study shows that it is possible to discover the signal of $W_H$ of the topflavor seesaw model in the semileptonic mode.

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