Neutrino Jets from High-Mass $W_R$ Gauge Bosons in TeV-Scale Left-Right Symmetric Models

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We re-examine the discovery potential at hadron colliders of high-mass right-handed (RH) gauge bosons $W_R$ - an inherent ingredient of Left-Right Symmetric Models (LRSM). We focus on the regime where the $W_R$ is very heavy compared to the heavy Majorana neutrino $N$, and investigate an alternative signature for $W_R \to N$ decays. The produced neutrinos are highly boosted in this mass regime. Subsequently, their decays via off-shell $W_R$ bosons to jets, i.e., $N \to \ell^\pm j j$ are highly collimated, forming a single neutrino jet ($j_N$). The final-state collider signature is then $\ell^\pm j N$, instead of the widely studied $\ell^\pm \ell^\pm j j$. Present search strategies are not sensitive to this hierarchical mass regime due to the breakdown of the collider signature definition. We take into account QCD corrections beyond next-to-leading order (NLO) that are important for high-mass Drell-Yan processes at the 13 TeV Large Hadron Collider (LHC). For the first time, we evaluate $W_R$ production at NLO with threshold resummation at next-to-next-to-leading logarithm (NNLL) matched to the threshold-improved parton distributions. With these improvements, we find that a $W_R$ of mass $M_{W_R} = 3 (4) [5]$ TeV and mass ratio of $(m_N/M_{W_R}) < 0.1$ can be discovered with a $5 - 6\sigma$ statistical significance at 13 TeV after 10 (100) [2000] fb$^{-1}$ of data. Extending the analysis to the hypothetical 100 TeV Very Large Hadron Collider (VLHC), $5\sigma$ can be obtained for $W_R$ masses up to $M_{W_R} = 15 (30)$ with approximately 100 fb$^{-1}$ (10 ab$^{-1}$). Conversely, with 0.9 (10) [150] fb$^{-1}$ of 13 TeV data, $M_{W_R} < 3 (4) [5]$ TeV and $(m_N/M_{W_R}) < 0.1$ can be excluded at 95% CL; with 100 fb$^{-1}$ (2.5 ab$^{-1}$) of 100 TeV data, $M_{W_R} < 22 (33)$ TeV can be excluded.

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I. INTRODUCTION

The observation of nonzero neutrino masses $m_{\nu}$ that have hierarchically smaller masses than all other elementary fermions in the Standard Model of Particle Physics (SM), and their non-trivial mixing provide unambiguous experimental evidence of physics beyond the SM (BSM). The natural explanation for such tiny masses is the so-called Seesaw Mechanism, where eV neutrino masses are generated from the $(B - L)$-violating operators at dimension-5 \[1, 2\]. At tree level, these operators can be generated by extending minimally \[3\] the SM field contents by right-handed (RH) neutrinos $N_R$ (Type I) \[4–9\], scalar SU(2)$_L$ triplets $\Delta_L$ (Type II) \[10–13\], or fermionic SU(2)$_L$ triplets $\Sigma$ (Type III) \[14\]. If kinematically accessible, these states can be observed at the 13 TeV Large Hadron Collider (LHC) or a hypothetical 100 TeV Very Large Hadron Collider (VLHC) \[15, 16\], thus giving conclusive evidence of the mass generation mechanism. For reviews of TeV-scale Seesaw models and their phenomenology, see Refs. \[17\].

An appealing renormalizable framework in which both Types I and II Seesaws can be embedded is the Left-Right Symmetry Model (LRSM) \[18–21\]. This is based on the gauge group

\[ \text{SU}(3)_c \otimes \text{SU}(2)_L \otimes \text{SU}(2)_R \otimes \text{U}(1)_{B-L}. \]

and postulates the restoration of parity symmetry at high energies. In addition to the SM particle content, the model consists of three generations of $N_R$, one $\Delta_L$, and an SU(2)$_R$ triplet scalar $\Delta_R$, all with non-trivial charges under the $B - L$ symmetry. After $\Delta_R$ acquires a vev $v_R$, much larger than the electroweak (EW) scale, $v_{\text{SM}} \approx 246$ GeV, the $SU(2)_R \times U(1)_{B-L}$ symmetry breaks down to the $U(1)'_{\nu}$ of the SM. Subsequently, the neutrinos $N_R$ and gauge bosons $W_R$ and $Z_R$ acquire masses $M_R$, $M_{W_R}$, and $M_{Z_R}$, respectively, that are proportional to $v_R$. While the masses of the gauge bosons depend on the weak gauge coupling $g_R = g$, the masses of $N_R$ are dependent on the Yukawa coupling $f_R$ of the $\Delta_R$ and lepton doublet interaction. The RH neutrino also interacts with the SM neutrino via its Yukawa interaction, generating Dirac masses $M_D$ after EW symmetry breaking (EWSB). For the Majorana mass $M_R$ much
heavier than the Dirac mass $M_D$, a Type I Seesaw is triggered, giving rise to Majorana masses for light neutrinos $\nu_m$ with $m_{\nu} \sim M^2_D/M_R$ and heavy neutrinos $N$ with $m_N \sim M_R$. As no symmetry relates the RH gauge and triplet Yukawa couplings, the $W_R$ and heavy neutrino may have widely separated masses, and offers a wider parameter space to test the LRSM.

The LRSM model can be tested either indirectly, through low energy experiments, or directly, through searches at high energy colliders (and references therein). In this work, we focus on the direct detection of the $W_R$ and $N$. For $M_{W_R} > m_N$, the hallmark hadron collider test of the LRSM is the spectacular lepton number ($L$) violating process

$$q_1 q_2 \rightarrow W_R^\pm \rightarrow N_R \ell^\pm \rightarrow \ell_1^\pm \ell_2^\pm W_R^{\mp} \rightarrow \ell_1^\pm \ell_2^\pm q_1 \bar{q}_2.$$  \hspace{1cm} (2)

The process, shown in Fig. 1, has been studied extensively. Searches by the ATLAS [55] and CMS [54] collaborations have excluded regions of the $(M_{W_R}, m_N)$ parameter space for $M_{W_R}$ up to several TeV (hundred GeV) [57]. However, for hierarchical masses, i.e., $(m_N/M_{W_R}) < 0.1$, the present search strategy is no longer sensitive. Complimentary dijet searches have similarly excluded $M_{W_R}$ below 2.5 - 3.5 TeV [59, 61].

In light of such stringent bounds, we re-examine the discovery potential of high-mass $W_R$ at hadron colliders. We focus on the situation where $N$ are hierarchically lighter than $W_R$, i.e., $(m_N/M_{W_R}) < 0.1$. In the process $p p \rightarrow W_R \rightarrow N \ell$, this leads to boosted $N$ with transverse momentum $p_T^N \sim M_{W_R}/2$. The decay products of $N$, which proceed dominantly through off-shell $W_R$ to quarks, $N \rightarrow W_R^\ell \rightarrow q \bar{q} \ell$, are subsequently collimated with parton separations scaling as $\Delta R_{ij} \sim 2 m_N/p_T^N \sim 4 m_N/M_{W_R}$. Hence, for $m_N/M_{W_R} \lesssim 0.1$, one has $\Delta R_{ij} \lesssim 0.4$, which falls below the electron isolation threshold in standard high-$p_T$ lepton searches at the 13 TeV LHC [62]. Indeed, the LHC sensitivity of Eq. (2) for such $(M_{W_R}, m_N)$ is considerably weaker, particularly in the electron channel [55].

This deficiency has been noted before, e.g., Refs. [38, 51, 63], but never explored in substantial detail.

After hadronization, the decay products of $N$ do not appear as individual, isolated objects, but instead as a single neutrino jet $j_N$. This is akin to the formation of top jets from boosted top quarks [64–68]. Thus, for $m_N \ll M_{W_R}$, $W_R - N$ production and decay appear in $pp$ collisions as the distinctive $p p \rightarrow W_R^\pm \rightarrow \ell^\pm j_N$.  \hspace{1cm} (3)

Despite the inclusive channel’s simplified topology, and hence larger SM backgrounds, it inherits much of the strong discriminating power of Eq. (2), including a fully reconstructible final state and no missing $p_T$ (MET), other than the hadronization and detector effects. We consider a search strategy for $W_R - N$ production and decay when $M_{W_R} > 3$ TeV and $(m_N/M_{W_R}) < 0.1$, while using a simple set of kinematical cuts on the effective two-body final state. We explore the discovery potential of observing Eq. (3) for the c.m. energies $\sqrt{s} = 13$ and 100 TeV, relevant for the LHC and VLHC.

Furthermore, determining if the $W_R$ gauge coupling $g_R$ equals the SM weak coupling $g$, a postulate of Eq. (1), requires precision knowledge of $W_R$ production rates. However, for such large $W_R$ masses, QCD corrections beyond next-to-leading order (NLO) are important at 13 TeV because of soft gluon radiation off initial-state partons. In light of this, we also calculate, for the first time, $W_R$ production at NLO with threshold resummation at next-to-next-to-leading logarithm (NNLL) matched to threshold-improved parton distributions functions (PDFs) [65, 70]. Previous predictions [61, 71, 72] have considered threshold resummation up to next-to-leading logarithm (NLL) [61, 62], but never matched to resummed PDFs. NLO+NNLL contributions improve the Born (NLO)-level predictions for $M_{W_R} = 4 - 5$ TeV by 40 - 140 (4 - 72)% at 13 TeV LHC.

With these improvements, we find that a $W_R$ of mass $M_{W_R} = 3$ (4) [5] TeV and $(m_N/M_{W_R}) < 0.1$ can be discovered with a 5 - 6σ statistical significance at 13 TeV after 10 (100) [2000] fb$^{-1}$. At 100 TeV with 0.1 (10) ab$^{-1}$, the 5σ reach extends to $M_{W_R} = 15$ (30) TeV. Conversely, with 0.9 (10) [150] fb$^{-1}$ of 13 TeV data, $M_{W_R} < 3$ (4) [5] TeV and $(m_N/M_{W_R}) < 0.1$ can be excluded at 95% CL; with 100 fb$^{-1}$ (2.5 ab$^{-1}$) of 100 TeV data, $M_{W_R} < 22$ (33) TeV can be excluded.

Our report is organized as follows: In Sec. II we briefly review the minimal LRSM (MLRSM) and current constraints. In Sec. III we present predictions up to NLO+NNLL for $W_R - N$ production and decay rates at hadron colliders. We explore the phenomenology of boosted heavy neutrinos and present our signal-versus-background analysis in Sec. IV. In Sec. V we summarize and conclude. We relegate technical details of our resummation calculation to App. A and implementation of the LRSM model files by Ref. [73] to App. B.

II. MINIMAL LEFT-RIGHT SYMMETRIC MODEL

Here, we briefly review main aspects of the MLRSM relevant to our study. For an expanded discussion, see, e.g., Refs. [79]. In Secs. II A and II B we address the masses of $W_R$ and $N$. In Sec. II C experimental constraints are
FIG. 1. Born diagram of $W_R$ production in hadron collisions and decay via $N$ to leptons and quarks. Figures are drawn using JaxoDraw [54].

reviewed. We reserve discussing the model’s scalar potential and its implementation into publicly available simulation model files [73] for App. B. As we use the files of Ref. [73], we adopt their notation.

The MLRSM [18–20] is based on the extended gauge group

$$SU(3)_c \otimes SU(2)_L \otimes SU(2)_R \otimes U(1)_{B-L}. \quad (4)$$

In addition to the SM fermion field content, there are three generations of RH neutrinos $N_R$. Quark and lepton multiplets are assigned the following gauge group representations:

$$Q_{L,i} = \left( \begin{array}{c} u_L \\ d_L \end{array} \right)_i : (3, 2, 1, \frac{2}{3}), \quad Q_{R,i} = \left( \begin{array}{c} u_R \\ d_R \end{array} \right)_i : (3, 1, 2, \frac{1}{3}),$$

$$\psi_{L,i} = \left( \begin{array}{c} \nu_L \\ e_L \end{array} \right)_i : (1, 2, 1, -1), \quad \psi_{R,i} = \left( \begin{array}{c} N_R \\ e_R \end{array} \right)_i : (1, 1, 2, -1). \quad (5)$$

In the above, $i = 1, \ldots, 3$, is the family index. $(B-L)$ charges are normalized such that the electric charge is given by

$$Q = I_3^L + I_3^R + (B-L)/2,$$

and $I_3^L(R)$ being the third isospin components of $SU(2)_L(R)$. The scalar sector consists of the following multiplets:

$$\Phi = \left( \begin{array}{c} \phi_0^+ \\ \phi_1^- \phi_2^+ \end{array} \right) : (1, 2, 2, 0), \quad (6)$$

$$\Delta_L = \left( \begin{array}{cc} \Delta_L^+/\sqrt{2} & \Delta_L^{++} \\ \Delta_L^+/\sqrt{2} & -\Delta_L^{--} \end{array} \right) : (1, 3, 1, 2), \quad \Delta_R = \left( \begin{array}{cc} \Delta_R^+/\sqrt{2} & \Delta_R^{++} \\ \Delta_R^+/\sqrt{2} & -\Delta_R^{--} \end{array} \right) : (1, 1, 3, 2).$$

At a scale much higher than the EW scale, $\Delta_R$ acquires a vev $v_R = \sqrt{2} \langle \Delta_R \rangle$. This triggers spontaneous breaking of the $LR$- and $(B-L)$-symmetries, and reduces Eq. (4) to the SM gauge group, i.e., $SU(2)_R \times U(1)_{B-L} \rightarrow U(1)_Y$.

The bidoublet $\Phi$ is responsible for Dirac masses and EW symmetry breaking (EWSB) after it acquires the vevs $\langle \Phi \rangle = \text{diag}(k_1, k_2)/\sqrt{2}$, where

$$k_{\pm} = k_1^2 \pm k_2^2 \quad \text{and} \quad k_+ = v_{\text{SM}} \approx 246 \text{ GeV}. \quad (7)$$

In the absence of fine-tuning, $k_1$, $k_2$ naturally scale as

$$\frac{k_2}{k_1} \sim \frac{m_b}{m_t} \ll 1. \quad (8)$$

$\Delta_L$ can also acquire a vev $v_L = \sqrt{2} \langle \Delta_L \rangle$. However, precision measurements of the $\rho/T$-parameter indicate $v_L$ is much smaller than the EW scale [22, 23]. For simplicity, we take $v_R$ and $k_{1,2}$ to be real, i.e., no CP violation, and $v_L = 0$. 

\[\text{FIG. 1. Born diagram of } W_R \text{ production in hadron collisions and decay via } N \text{ to leptons and quarks. Figures are drawn using JaxoDraw [54].}\]
A. Charged Gauge Boson Masses

After LR and EWSB, the charged vector boson (squared) mass matrix in the gauge, i.e., \((W_L, W_R)\), basis is given by

$$
\mathcal{M}_W = \frac{g^2}{4} \begin{pmatrix}
    k_1^2 + k_2^2 + 2v_L^2 & 2k_1k_2 \\
    2k_1k_2 & k_1^2 + k_2^2 + 2v_R^2
\end{pmatrix}.
$$

The gauge states are related to the mass eigenstates, i.e., \((W_1, W_2)\) with \(M_{W_2} > M_{W_1}\), by

$$
\begin{pmatrix}
    W_1 \\
    W_2
\end{pmatrix} = \begin{pmatrix}
    \cos \xi & \sin \xi \\
    -\sin \xi & \cos \xi
\end{pmatrix} \begin{pmatrix}
    W_L \\
    W_R
\end{pmatrix},
$$

where the \(W_L - W_R\) mixing parameter \(\xi\) is

$$
\tan 2\xi = \frac{2k_1k_2}{v_R^2 - v_L^2}.
$$

Under the vev hierarchy

$$
v_R \gg k_+ \geq k_1 \geq k_- \gg k_2 \gg v_L \sim 0,
$$

the vector boson masses simplify to

$$
M_{W_1} \approx M_{W_L} = \frac{g}{2} k_+ \quad \text{and} \quad M_{W_2} \approx M_{W_R} = \frac{g}{\sqrt{2}} v_R,
$$

implying that the \(W_1\) (\(W_2\)) mass state is closely aligned with the \(W_L\) (\(W_R\)) gauge state. Hence, for the remainder of the text, we refer to \(W_1\) (\(W_2\)) as \(W_L\) (\(W_R\)).

B. Neutrino Masses

The leptonic Yukawa couplings for generations \(i\) and \(j\) are given by

$$
\mathcal{L}_Y = -h_{ij}\bar{\psi}_L \Phi \psi_{R_j} - h_{ij}\bar{\psi}_L \tilde{\Phi} \psi_{R_j} - f_{L_{ji}} \bar{\psi}^T_L C \sigma_2 \Delta_L \psi_{L_j} - f_{R_{ji}} \bar{\psi}^T_R C \sigma_2 \Delta_R \psi_{R_j} + \text{H.c.},
$$

where \(C\) denotes the charge conjugation operator and \(\tilde{\Phi} = \sigma_2 \Phi^* \sigma_2\). After LR and EWSB, RH Majorana, LH Majorana, and Dirac neutrino mass matrices, respectively, of the form

$$
M_R = \sqrt{2} v_R f_R, \quad M_L = \sqrt{2} v_L f_L, \quad M_D = \frac{1}{\sqrt{2}} \left( k_1 h + k_2 \tilde{h} \right),
$$

are spontaneously generated. The \(3 \times 3\) matrices in Eq. (15) can be combined such that in the gauge basis, i.e., \((\nu_L, \nu_R, N_{R1}, \ldots)\), the \(6 \times 6\) neutrino mass matrix is given by

$$
\mathcal{M}_\nu = \begin{pmatrix}
    M_L & M_D \\
    M_D^T & M_R
\end{pmatrix},
$$

and can be diagonalized via the unitary matrix \(\tilde{V}\):

$$
\mathcal{M}_\nu^{diag} = \tilde{V}^T \mathcal{M}_\nu \tilde{V} = \begin{pmatrix}
    M^{diag}_\nu & 0 \\
    0 & M^{diag}_N
\end{pmatrix}.
$$

\(M^{diag}_\nu = \text{diag}(m_1, m_2, m_3)\) and \(M^{diag}_N = \text{diag}(m_{N_1}, m_{N_2}, m_{N_3})\) are the light neutrino and heavy neutrino masses, respectively. For the vev hierarchy of Eq. (12), \(\tilde{V}\) is \([74, 75]\)

$$
\tilde{V} = \begin{pmatrix}
    (1 + \zeta^* \zeta^T)^{-1/2} & \zeta^*(1 + \zeta^* \zeta^T)^{-1/2} \\
    -\zeta^T(1 + \zeta^* \zeta^T)^{-1/2} & (1 + \zeta^* \zeta^T)^{-1/2}
\end{pmatrix} \begin{pmatrix}
    U_L & 0 \\
    0 & Y_R
\end{pmatrix} = \begin{pmatrix}
    U & V \\
    X & Y
\end{pmatrix},
$$

where \(\zeta^* = M_D M_R^{-1}\) and \(U_L, Y_R\) are unitary matrices that diagonalize \(\tilde{M}_\nu\) and \(\tilde{M}_N\):

$$
M^{diag}_\nu = U_L^T \tilde{M}_\nu U_L \quad \text{and} \quad M^{diag}_N = Y_R^T \tilde{M}_N Y_R.
$$
\( \tilde{M}_\nu \) and \( \tilde{M}_N \) are related to the mass matrices in Eq. (15) by the Seesaw relations \( [4, 13] \)

\[
\tilde{M}_\nu \simeq M_e - M_D M_R^{-1} M_D^T \quad \text{and} \quad \tilde{M}_N \simeq M_R.
\]  

(20)

In the notation of Refs. [41, 76], after rotating the charged leptons from the flavor basis into the mass basis, which for simplicity we take to be a trivial rotation, the \( U_{\nu_{\ell}} (\nu_{\ell{N}_{null}}) \) of Eq. (18) denotes the large, \( \mathcal{O}(1) \) mixing between the LH (RH) lepton flavor state \( \ell \) (\( \ell = e, \mu, \tau \)) and light (heavy) neutrino mass eigenstate \( \nu_n \) (\( \nu_{null} \)). Similarly, \( V_{\ell N_{null}} (X_{\nu_{null}}) \) denotes the suppressed, \( \mathcal{O}(m_\ell/m_N) \) mixing between the LH (RH) lepton flavor state \( \ell \) and heavy (light) neutrino mass eigenstate \( N_{null} \) (\( \nu_{null} \)).

C. Experimental Constraints

Here, we review the most stringent constraints on the MLRSM.

1. Collider Bounds on \((M_{W_R}, m_N)\) from \( \ell^\pm \ell^\pm jj \) searches: Searches by the ATLAS experiment for \( pp \to e^\pm e^\pm jj \) (\( \mu^\pm \mu^\pm jj \)) mediated by \( W_R \) and \( N \) excludes at \( \sqrt{s} = 8 \) TeV \([55] \):

\[
M_{W_R} < 1.5 \ (2.7) \ \text{TeV} \quad \text{at} \quad 95\% \ \text{C.L.} \quad \text{with} \quad \mathcal{L} = 20.3 \ \text{fb}^{-1}.
\]

(21)

The sensitivity disparity is due a failing isolated electron-jet criterion when \( m_N/M_{W_R} \lesssim 0.1 \) \([55]\) and is the point of our study. Limits from CMS are comparable \([56]\).

2. Collider Bounds on \( M_{W_R} \) from dijet searches: Searches by the ATLAS (CMS) experiment for a sequential SM \( W' \to jj \), excludes at \( \sqrt{s} = 13 \) TeV \([54, 61] \):

\[
M_{W_{SM}} < 2.6 \ (2.6) \ \text{TeV} \quad \text{at} \quad 95\% \ \text{C.L.} \quad \text{with} \quad \mathcal{L} = 3.6 \ (2.4) \ \text{fb}^{-1}.
\]

(22)

3. Limits on \( W_R \) and Higgs masses from neutral hadron transitions: Analyses of \( \Delta F = 2 \) transitions in neutral \( K \) and \( B_{d,s} \) systems and neutron EDM assuming generalized charge (parity) in the MLRSM exclude \([31, 33]\):

\[
M_{W_R} < 2.9 - 20 \ \text{TeV} \quad \text{at} \quad 95\% \ \text{C.L.}
\]

\[
m_{\text{FCNH}} < 20 \ \text{TeV} \quad \text{at} \quad 95\% \ \text{C.L.}
\]

(23)

(24)

where the range over \( M_{W_R} \) is based on theoretical arguments and \( m_{\text{FCNH}} \) is the mass of the lightest Higgs mediating flavor changing neutral transitions.

4. Searches for \( 0\nu\beta\beta \): In MLRSM, the gauge boson \( W_R \) together with \( N_i \) can give a saturating contribution in \( 0\nu\beta\beta \). Non-observation of this LNV process hence constrains the masses of \( W_R \) and \( N_i \) as \( \sum_i \frac{V_{ij}^2}{M_{W_R}} \lesssim (0.082 - 0.076) \) TeV\(^{-5} \), using the 90\% C.L half-life limit from KamLAND-Zen \( T_{1/2}^{0\nu} \geq 1.07 \times 10^{26} \) yrs \([77]\). For a \( M_{W_R} \) of 3 TeV (5 TeV) this implies a lower limit on the \( m_N \geq 150 - 162 \) GeV (19.5 - 21 GeV) \([31]\).

III. PROPERTIES OF \( W_R \) AND \( N \) AT HADRON COLLIDERS

In this section, we present production and decay rates of \( W_R \) and \( N \) to leptons and jets, with \( m_N \ll M_{W_R} \), at the 13 TeV LHC and 100 TeV VLHC.

In the MLRSM, the \( W_R \) interaction with quarks is given by the Lagrangian

\[
\mathcal{L}_{W_R-q-q'} = -\frac{g}{\sqrt{2}} \sum_{i,j=u,d,...} \bar{u}_i V_{ij}^{\text{CKM}} W_{R_R}^{+} q_i \gamma^\mu P_R d_j + \text{H.c.,}
\]

(25)

where \( u_i(d_j) \) is an up-(down-)type quark of flavor \( i(j) \); \( V_{ij}^{\text{CKM}} \) is the RH Cabbibo-Kobayashi-Masakawa (CKM) matrix; and \( P_R(L) = \frac{1}{2}(1 \pm \gamma^5) \) denotes the RH(LH) chiral projection operator. For either generalized charge conjugation or parity, one expects \( |V_{ij}^{\text{CKM}}| \approx |V_{ij}^{\text{CKM}}| \), up to \( \mathcal{O}(m_q/m_H) \) contributions for the latter case \([78, 82]\). Hence, we can assume, for simplicity, four massless quarks and take \( V_{ij}^{\text{CKM}} \) to be diagonal with unit entries.
The $W_R$ coupling to six heavy ($N_m$) and light ($\nu_m$) neutrinos is parametrized by \[41, 70\]

$$
\mathcal{L}_{W_R \ell - \nu/N} = -\frac{g}{\sqrt{2}} \sum_{\ell=e,\mu,\tau} \left[ \sum_{m=1}^{3} \bar{\nu}_m X_{\ell m} + \sum_{m'=4}^{6} N_{m'} Y_{\ell m'} \right] W_{R\mu}^\pm \gamma^\mu P_R \ell^- + \text{H.c.,}
$$

(26)

where mixing matrices $X_{\ell m}$ and $Y_{\ell m'}$ are defined in Sec. II B. We consider only the lightest heavy neutrino, denoted simply by $N$, and neglect heavier mass eigenstates. For simplicity, we assume diagonal neutrino mixing with maximal coupling to electron-flavor leptons:

$$
|Y_{eN}| = 1, \quad |Y_{\mu N}| = |Y_{\tau N}| = |X_{\ell m}| = 0.
$$

(27)

Choosing instead maximal coupling to muons, i.e., $|Y_{\mu N}| = 1$, or large $e-\mu$ mixing, i.e., $|Y_{eN}| \sim |Y_{\mu N}|$, has little impact on our analysis due to the long lifetime of the muon. On the other hand, the $\tau\ell$ final state requires specialized cuts to account for $\tau$ decays to light neutrinos. For more details, see Sec. IV A.

For numerical results, SM inputs are taken from the 2014 Particle Data Group \[83\]:

$$
\alpha_{\text{MS}}(M_Z) = 1/127.940, \quad M_Z = 91.1876 \text{ GeV}, \quad \sin^2\theta_W(M_Z) = 0.23126.
$$

(28)

PDFs and $\alpha_s(\mu_f)$ are extracted using the LHAPDF 6.1.6 libraries \[84\]. The factorization ($\mu_f$) and renormalization ($\mu_f$) scales are set to $\mu_0 = M_W/r_n$ everywhere. For LO- and NLO-accurate calculations, we use the NNPDF 3.0 NLO pdf set \[85\]. For NLO+NNLL calculations, we use the threshold-improved NNPDF 3.0 NNLO+NNLL PDF set \[86\]. This choice follows from the unavailability of an NLO+NNLL PDF set and our desire to ascertain the effects of resummation at NNLL. Formally, the induced uncertainty from our PDF mismatching in the LO and NLO+NNLL calculations is $O(\alpha_s)$ and $O(\alpha_s^2)$, respectively, and beyond our claimed accuracy. Numerically, this leads to LO cross sections that are 10% smaller than those calculated with LO PDFs.

For additional computational details, see App. A.

### A. $W_R$ Production at NLO+NNLL

At Fixed Order (FO) accuracy, we calculate the inclusive production cross section for

$$
pp \to W_R^\pm + X,
$$

(29)

where $X$ is anything, via the usual application of the Collinear Factorization Theorem:

$$
\sigma^{\text{FO}}(pp \to W_R + X) = \sum_{a,b=q,q',g} \int_{\tau_0}^{1} d\tau \mathcal{L}_{ab}(\tau, \mu_f) \delta^{\text{FO}}(ab \to W_R), \quad \tau_0 = \frac{M_{W_R}^2}{s}.
$$

(30)

The luminosity $\mathcal{L}(\tau)$ of parton pair $ab$ in $pp$ collisions given by

$$
\mathcal{L}_{ab}(\tau, \mu_f) = \frac{1}{1 + \delta_{ab}} \int_{\tau}^{1} d\xi \frac{1}{\xi} \left[ f_{a/p}(\xi_1, \mu_f) f_{b/p}(\xi_2, \mu_f) + f_{a/p}(\xi_2, \mu_f) f_{b/p}(\xi_1, \mu_f) \right],
$$

(31)

$$
\xi_2 = \frac{\tau}{\xi_1}.
$$

(32)

The PDFs $f_{a/p}(\xi, \mu_f)$ represent the likelihood of observing parton $a$ in proton $p$ possessing a longitudinal momentum fraction $\xi = E_a / E_p = p_z / p_T$, and (re)sum arbitrary collinear parton emissions up to a factorization scale $\mu_f$. The partonic c.m. energy $\sqrt{s}$ is related to the hadronic (beam) c.m. energy $\sqrt{s}$ by the hadronic threshold variable

$$
\tau = \xi_1 \xi_2 = \frac{\hat{s}}{s}, \quad \tau_0 \leq \tau < 1,
$$

(33)

and extends to the kinematic threshold $\tau_0$, below which Eq. (29) is kinematically forbidden.

Partonic scattering rates $\hat{s}$ are evaluated via helicity amplitudes, and use the CUBA libraries \[86\] to handle Monte Carlo integration. NLO in QCD corrections are obtained using the Phase Space Slicing method \[87\]–\[90\] and exploit factorization properties of Drell-Yan (DY) currents; see appendices of Refs. \[90\]–\[91\]. LO and NLO results are checked against literature \[61\]–\[73\] and MG5_aMC@NLO v2.3.3 (MG5) \[92\] assuming $M_{W_R} = M_W$. 
Beyond FO, Eq. (30) can be generalized to include the arbitrary, initial-state emission of soft gluons, i.e., with energies much smaller than the hard scattering process scale $Q$. The interpretation of $\hat{\sigma}$ also generalizes to include both the hard process, $q \bar{q}' \to W_R^\pm$ with $Q = M_{W_R}$, and the factorized soft radiation off the $q, q'$ initial states. Schematically, the definitions of the hadronic, partonic, and hard components for the inclusive production of a generic color-singlet boson $B$ are drawn in Fig. 2. Necessarily, the inequality $s > \hat{s} \geq Q^2$ holds.

Soft radiation becomes important when the hard scale approaches the partonic scale, i.e., when the partonic threshold variable $z$ approaches one:

$$z \equiv \frac{Q^2}{\hat{s}} = \frac{M^2_{W_R}}{\hat{s}} = \frac{\tau_0}{\tau} \to 1.$$  (35)

In this kinematic regime, which can be satisfied at $Q^2 \ll s$ as in Higgs production via GF or when $Q^2 \sim s$ as in the present case of high-mass DY, soft radiation give rise to numerically large logarithms that require resummation in order to restore perturbativity of Eq. (30).

To carry out the resummation, we follow the procedure (and largely notations) of Refs. [96–98], and write a generalized form of Eq. (30) in terms of $\tau, z$, and $\tau_0$:

$$\sigma_{FO}(pp \to W_R + X) = \sum_{a,b=q,q',g} \int_{\tau_0}^{1} d\tau \int_{0}^{1} dz \, \delta \left( z - \frac{\tau_0}{\tau} \right) L_{ab}(\tau) \hat{\sigma}_{ab}^{FO}(ab \to W_R).$$  (36)

For inclusive $W_R$ production, $\hat{\sigma}_{ab}^{FO}$ can be expressed as

$$\hat{\sigma}_{ab}^{FO} \equiv \hat{\sigma}^{FO}(ab \to W_R) = \sigma_0 \times z \times \Delta_{ab}^{FO}(z).$$  (37)

The constant term $\sigma_0$ for gauge coupling $g_R^2 = g^2 = 4\pi\alpha / \sin^2 \theta_W$ is

$$\sigma_0 = \frac{g^2_R \pi |V_{CKM}|^2}{4N_c M^2_{W_R}},$$  (38)

and is related to the usual LO partonic formula by

$$\hat{\sigma}^{LO}(ab \to W_R) = \sigma_0 \times M^2_{W_R} \times \delta(\hat{s} - M^2_{W_R}) = \sigma_0 \times z \times \delta(1-z).$$  (39)

Hence, one may identify up to $O(\alpha_s)$, $\Delta_{ab}^{FO}(z) \approx \delta(1-z) + O(\alpha_s)$.

If working with pQCD, the threshold resummed cross section can be efficiently obtained after writing the hadronic cross section in so-called Mellin-space. For the function $h(x)$, the $N$th-moment of its Mellin transform and inverse...
FIG. 3. Upper panel: As a function of $M_{W_R}$, $pp \to W_R$ production cross section for $\sqrt{s} =$ (a) 13 and (b) 100 TeV, at LO (solid), NLO (dash), and NLO+NNLL (dash-dot) with 1σ PDF uncertainty (shaded), as well as $\sigma^{NLO+NNLL}(pp \to W_R) \times \mathcal{B}(W_R \to N e) \times \mathcal{B}(N \to eqq')$ (dot). Lower: NLO (dash) and NLO+NNLL (dash-dot) $K$-factors and PDF uncertainties.

Mellin transform with respect to $x$ are,

$$h_N \equiv \mathcal{M}[h(x); N] = \int_0^1 dx \ x^{N-1} h(x),$$

$$h(x) = \mathcal{M}^{-1}[h_N; x] = \frac{1}{2\pi i} \int_{c-i\infty}^{c+i\infty} dN \ x^{-N} h_N,$$

where $c \in \mathbb{R}$ is to the right of all singularities in $h_N$. The Mellin transform of Eq. (36) at LO with respect to $\tau_0$, gives

$$\sigma_N^{LO} = \int_0^1 d\tau_0 \ \tau_0^{N-1} \times \sigma_0 \mathcal{L}_{qq'}^{(N+1)} \times \Delta_{qq'}^{LO}(N+1),$$

revealing an explicit factorization into a product of the luminosity and soft coefficient, normalized by the Born weight $\sigma_0$. We drop the summation over $a,b = g$ as the $gq$, $gq'$, and $gg$ initial states do not contribute to $W_R$ production at LO.

The advantage of working in Mellin-space is this explicit factorization. Exploiting that in the soft limit gauge radiation amplitudes reduce to their color-connected Born amplitudes, resummation reduces to the simple procedure of replacing the LO soft coefficient $\Delta_{qq'}^{LO}$ with its resummed analogue $\Delta_{qq'}^{Res}$ \cite{93,94}. Thus, the threshold-resummed $pp \to W_R$ cross section in Mellin-space is

$$\sigma_N^{Res} = \sigma_0 \mathcal{L}_{qq'}^{(N+1)} \times \Delta_{qq'}^{Res}(N+1),$$

and in momentum space by Mellin inverse of the above with respect to $\tau_0$:

$$\sigma^{Res}(pp \to W_R + X) = \frac{\sigma_0}{2\pi i} \int_{c-i\infty}^{c+i\infty} dN \ \tau_0^{-N} \times \mathcal{L}_{qq'}^{(N+1)} \times \Delta_{qq'}^{Res}(N+1).$$

We approximate the luminosity function $\mathcal{L}(\tau)$ using the Chebyshev polynomial approximation \cite{95,96}, which can be Mellin-transformed analytically, and choose the integration path according to the Minimal Prescription (MP) procedure \cite{96}. See App. A for more details.
Matching resummed and FO calculations beyond LO requires subtracting the soft contributions common to both calculations to avoid phase space double counting. For a FO result at N^4LO, this can be done by Taylor-expanding \( \sigma^{\text{Res.}} \) up to \( \mathcal{O}(\alpha_s^6) \), subtracting these terms from \( \sigma^{\text{Res.}} \), and adding the N^kLO calculation to the residual resummed expression. One may interpret this procedure as augmenting with approximate \( \mathcal{O}(\alpha_s^6) \) terms in \( \sigma^{\text{Res.}} \), i.e., soft/non-hard, with the full \( \mathcal{O}(\alpha_s^6) \) calculation, which describes accurately both soft and hard radiation. Subsequently, \( W_R \) production matched at N^4LO+N^NLL is given by

\[
\sigma^{N^4\text{LO}+\text{N}^\text{NLL}}(pp \to W_R + X) = \sigma^{N^4\text{LO}} + \sigma^{N^\text{NLL}} - \sum_{l=0}^{k} \frac{\alpha_l^l}{l!} \left[ \frac{d \sigma^{\text{N}^\text{NLL}}}{d \alpha_s^l} \right]_{\alpha_s=0}.
\]

(45)

In Fig. [3] we show the total inclusive \( pp \to W_R \) cross section at NLO+NNLL (dash-dot) with PDF uncertainty (shaded), NLO (dash), and LO (solid) at (a) 13 and (b) 100 TeV. The production rates at 13 (100) TeV span approximately:

\[
2 \text{ fb} - 40 \text{ pb} (90 \text{ ab} - 930 \text{ pb}) \quad \text{for} \quad M_{W_R} = 1 - 5 (1 - 33) \text{ TeV}.
\]

(46)

In the lower panel are the NLO+NNLL and NLO K-factors, defined respectively as

\[
K^{\text{NLO+NNLL}} = \frac{\sigma^{\text{NLO+NNLL}}}{\sigma^{\text{LO}}} \quad \text{and} \quad K^{\text{NLO}} = \frac{\sigma^{\text{NLO}}}{\sigma^{\text{LO}}}.
\]

(47)

The NLO+NNLL (dash-dot) and NLO (dash) K-factors with uncertainties span roughly:

\[
K^{\text{NLO+NNLL}} : 1.2 - 2.4 (1.2 - 1.5) \quad \text{and} \quad K^{\text{NLO}} : 1.2 - 1.4 (1.1 - 1.3).
\]

At 13 and 100 TeV, we observe that the effects of resummation become important with respect to the NLO rate at \( \tau_0 \approx 0.3 \). At 13 TeV, the resummed corrections for \( \tau_0 > 0.3 \) are very large, increasing the Born (NLO) predictions by 40 – 140 (4 – 70)\% for \( M_{W_R} = 4 - 5 \text{ TeV} \). The largeness of the 13 TeV K-factors for \( M_{W_R} \gtrsim 4 \text{ TeV} \) does not indicate the breakdown of perturbation theory. Rather, it demonstrates the importance of soft radiation as \( \tau_0 \to 1 \), and is typical for processes near the boundaries of phase space [60]. For the DY process, this is particularly important for \( \tau_0 \gtrsim 0.1 \) [100]. This is exemplified at 100 TeV by the reduced importance of resummation for comparable \( M_{W_R} \) (smaller \( \tau_0 \)). Despite the largeness of the PDF uncertainties at large \( M_{W_R} \), the NLO+NNLL central value remains within the NLO uncertainty, as seen in the lower panel of Fig. [3(a)]. See Sec. [III.D] for further discussions on uncertainties. Away from threshold, the resummed calculation converges to the FO result, consistent with expectations [60]. For select \( M_{W_R} \), we summarize our NLO and NLO+NNLL results in Tb. [4].

### B. \( W_R \) Decay

As discussed in Sec. [II.A] and in Sec. [II.B], \( W_R - W_L \) and \( N_i - \nu_i \) mixing are negligibly small and \( m_{\text{FCNH}} \gg M_{W_R} \). Subsequently, for \( m_N < M_{W_R} \), the only open \( W_R \) decay modes are to quark and \( \pm N \) pairs. The corresponding
The branching fraction of $A$ to final-state $X_i$ is defined as
\[ \text{BR}(A \to X_i) = \frac{\Gamma (A \to X_i)}{\sum_i \Gamma (A \to X_i)}. \] (55)

In the large $M_{WR}$ limit, the $W_R$ branching fractions converge to the asymptotic values,
\[ \text{BR}(W_R \to q\bar{q}) \approx 2 \times \text{BR}(W_R \to tb) \approx \frac{2N_c}{3N_c + 1} = 60\%, \] (56)
\[ \text{BR}(W_R \to Ne) \approx \frac{1}{3N_c + 1} = 10\%. \] (57)

In the upper (lower) panel of Fig. 4(a) we show the total $W_R$ decay width (branching fraction) for $M_{WR} > 1$ TeV and fixed $m_N/M_{WR}$ ratios of $\sqrt{r_N} = 0.01$ (dash), 0.1 (solid), 0.5 (dot), and 0.75 (dot-dash). Similar to the EW gauge bosons, the $W_R$ in this model has a narrow width for all values of $M_{WR}$, with $\Gamma_{WR}/M_{WR}$ scaling as:
\[ \frac{\Gamma_{WR}}{M_{WR}} \sim \frac{g^2}{48\pi} (3N_c + 1) \approx 2.8\%. \] (58)

This justifies the use of the Narrow Width Approximation (NWA). Furthermore, as $pp \to W_R$ is a DY process, its factorization properties imply that the NLO and NLO+NNLL corrections to its on-shell production and decay to $N$ are equivalent to the production-only corrections, i.e.,
\[ \sigma^{\text{NLO(}+\text{NNLL)}}(pp \to W_R \to e^+e^-N) \approx \sigma^{\text{NLO(}+\text{NNLL)}}(pp \to W_R) \times \text{BR}(W_R \to e^+e^-). \] (59)

In the lower panel of Fig. 4(a) we observe that the $W_R$ branching fractions remain virtually independent of $m_N$ and attain its maximum branching of $\text{BR}(W_R \to Ne^\pm) \approx 0.1$. For 13 (100) TeV and $(m_N/M_{WR}) = 0.1$ the $pp \to W_R \to Ne^\pm$ cross section [Eq. (59)] spans:
\[ 180 \text{ pb} - 15 \text{ fb} \ (100 \text{ pb} - 350 \text{ fb}) \quad \text{for} \quad M_{WR} = 3 - 5 \ (5 - 25) \text{ TeV}. \] (60)

For representative $(M_{WR}, m_N)$, we summarize our results in column 2 of Tb. III.
Following Eq. (27), such decays vanish at tree-level and, therefore, are not considered in the analysis. For neutrino mixing with SM neutrinos; the rate is controlled by the tiny mixing parameter $\sqrt{m_N/m_{W_R}}$. The validity of this approximation for neutrino decay to SM EW bosons is justified but suggests $N$ may be long-lived. Values of $\Gamma_N$ for representative $m_N$ used in this study are given in Tab. [II].

C. $N$ Decays

In our scenario, the heavy neutrino dominantly decays to the three-body final state

$$N \to e^\pm W^\mp_{R*} \to e^\pm q \bar{q}.$$  \hspace{1cm} (61)

Both $e^+$ and $e^-$ are allowed in the final state due to the Majorana nature of $N$. If kinematically accessible, the heavy neutrino can also decay to $t$ and 6 quarks, with the final state $e^\pm tb$. In principle, $N$ can also decay to SM EW bosons via mixing with SM neutrinos; the rate is controlled by the tiny mixing parameter $|X_{tN}|^2 \sim 1 - |Y_{\tau N}|^2 \sim O(m^2_\tau/m^2_N)$. Following Eq. (27), such decays vanish at tree-level and, therefore, are not considered in the analysis. For $m_N \ll M_{W_R}$, the partial widths of $N$ are

$$\Gamma_N = 2N_e |Y_{tN}|^2 |V_{CKM}^{tb}|^2 \frac{y_t^2}{3 \cdot 211 \cdot \pi^3 M_{W_R}^4} m_N."$$ (62)

$$\Gamma(N \to e^\pm t \bar{b}) = 2N_e |Y_{tN}|^2 |V_{CKM}^{tb}|^2 \frac{y_t^2}{3 \cdot 211 \cdot \pi^3 M_{W_R}^4} (1 - 8y_t^2 + 8y_t^3 - y_t^4 - 12y_t^2 \log y_t).$$ \hspace{1cm} (63)

The validity of this approximation for $m_N/M_{W_R} \sim 0.1$ has been checked against MG5. For our choice of mixing, the total $N$ width is

$$\Gamma_N = 2\Gamma(N \to e^\pm q \bar{q}) + \Gamma(N \to e^\pm t \bar{b})$$ \hspace{1cm} (64)

$$= 2N_e |Y_{tN}|^2 |V_{CKM}^{tb}|^2 \frac{y_t^2}{3 \cdot 211 \cdot \pi^3 M_{W_R}^4} [3 - 8y_t^2 + 8y_t^3 - y_t^4 - 12y_t^2 \log y_t].$$ \hspace{1cm} (65)

and implies that $\Gamma_N/m_N$ scales as

$$\frac{\Gamma_N}{m_N} = \frac{g^4}{2^{10} \pi^3} \left( \frac{m_N}{M_{W_R}} \right)^4 \sim 5 \cdot 10^{-6} \times \left( \frac{m_N}{M_{W_R}} \right)^4 \ll 1.$$ \hspace{1cm} (66)

Hence, application of the NWA in $N$ decays is justified but suggests $N$ may be long-lived. Values of $\Gamma_N$ for representative $M_{W_R}$ and $m_N$ used in this study are given in Tab. [II].
TABLE III. Cross section times branching ratio predictions for $pp \rightarrow W_R^\pm \rightarrow Ne^\pm$, with subsequent decay of $N$ to leptons and quarks, for select $(M_{W_R}, m_N)$.

In Fig. 4(b) we plot $\Gamma_N$ as a function of $m_N$ for representative $M_{W_R}$ and $(m_N/M_{W_R})$ ratios; in the lower panel we show the mean flight distances

$$d_0 = v\tau_0 = \beta\gamma hc/\Gamma_N, \quad \beta\gamma = \frac{(1 - r_N)}{2\sqrt{r_N}}$$

(67)

For $m_N = 30 - 1000$ GeV, we find

$$\frac{m_N}{M_{W_R}} = 0.1 \text{ (solid)} : \Gamma_N \sim 10^{-8} - 10^{-6} \text{ GeV},$$

(68)

$$\frac{m_N}{M_{W_R}} = 0.01 \text{ (dash)} : \Gamma_N \sim 10^{-12} - 10^{-10} \text{ GeV}.$$  

(69)

The corresponding mean flight distances span

$$\frac{m_N}{M_{W_R}} = 0.1 \text{ (solid)} : d_0 \sim 10^{-7} - 10^{-5} \text{ mm},$$

(70)

$$\frac{m_N}{M_{W_R}} = 0.01 \text{ (dash)} : d_0 \sim 10^{-2} - 3 \text{ mm}.$$  

(71)

This implies that for $N$ much lighter than $M_{W_R}$, i.e., $m_N/M_{W_R} < 0.01$, heavy neutrinos appear in detector experiments as displaced vertices, not prompt decays. However, such a scenario is not reasonable within the spirit of the LRSM model.

Supposing $m_N/M_{W_R} < 0.01$ and using expressions for $m_N, M_{W_R}$ in Sec. IIB and Sec. IIA the Yukawa couplings of the heavy neutrino $N$ to the triplet Higgs are restricted to $f_R < 3 \times 10^{-3}$. This is comparable to generation I and II quark SM Yukawa couplings. However, taking $m_N \sim \mathcal{O}(10)$ GeV, a (vanilla) Type I Seesaw then requires for light neutrino masses $m_{\nu_m} \sim 0.1$ eV a Dirac neutrino mass of $m_D \sim 30$ KeV, or a Yukawa coupling $\mathcal{O}(15 - 20) \times$ smaller than the SM electron Yukawa. Though not forbidden, this is contrary to the Seesaw spirit of explaining light neutrino masses without excessively small couplings.

From Eq. (65) the $N$ branching fractions are independent of $M_{W_R}$ and are given by

$$\text{BR}(N \rightarrow e^\pm q \overline{q}) = \begin{cases} 
1, & m_N \leq m_t, \\
2 \times \frac{m_N}{M_{W_R}} \left( 3 - 8y_t + 8y_t^2 - y_t^4 - 12y_t^2 \log y_t \right), & m_N > m_t.
\end{cases}$$

(72)

$$\text{BR}(N \rightarrow e^\pm t \bar{b}) = \begin{cases} 
1 - 8y_t + 8y_t^3 - y_t^4 - 12y_t^2 \log y_t, & m_N > m_t, \\
3 - 8y_t + 8y_t^3 - y_t^4 - 12y_t^2 \log y_t, & m_N > m_t.
\end{cases}$$

(73)

For $M_{W_R} \gg m_N \gg m_t$, one finds asymptotically

$$\text{BR}(N \rightarrow e^\pm q \overline{q}) \approx 2 \times \text{BR}(N \rightarrow e^\pm t \bar{b}) \approx \frac{2}{3}.$$  

(74)

Consequently, the 13 and 100 TeV cross sections for the process

$$pp \rightarrow W_R \rightarrow Ne \rightarrow ee q \overline{q}$$

(75)
in the NWA approximation can be given in terms of Eq. \(69\):

\[
\sigma^{NLO+NNLL}(pp \rightarrow W_R^+ \rightarrow Ne^+ e^- q\bar{q}) \approx \sigma^{NLO+NNLL}(pp \rightarrow W_R^+) \\
\times \text{BR}(W_R \rightarrow Ne) \\
\times \text{BR}(N \rightarrow e^+ q\bar{q})
\] (76)

The total production rate for Eq. \(76\) for representative \((M_{W_R}, m_N)\) are summarized in column 3 of Tb. II and for \(m_N/M_{W_R} = 0.1\) plotted in Fig. 3 (dot). We find that the total 13 (100) TeV rate spans approximately

\[10^{-1} - 4 \times 10^4 \, (10^{-3} - 10^5) \, \text{fb} \quad \text{for} \quad M_{W_R} = 1 - 5 \, (35) \, \text{TeV}.\] (77)

D. PDF and Scale Uncertainties

To estimate the impact of higher order terms in the QCD perturbative series that are not calculated in the \(W_R\) production cross section, we vary the factorization and renormalization scales about the default choice of \(\mu_0 = M_{W_R}\) up and down by a factor of two. We present results normalized to the cross section at the default scale. In the lower panel of each plot is the \(K\)-factor as defined in Eq. \(47\).

In Fig. 5 we show the effect of scale variation on the NLO cross section at (a) 13 and (b) 100 TeV for a range of \(W_R\) masses. At NLO, it can be seen at both 13 and 100 TeV that increasing (decreasing) the default scale lowers (raises) the total cross section, except for very low \(W_R\) masses at 100 TeV, a feature common to high-mass DY processes \cite{r101}. In addition, the \(K\)-factor also steadily increases with mass indicating the growing importance of higher order corrections in such scenarios. In both the 13 and 100 TeV cases, the scale variation results in a 2−5% uncertainty to the total cross section.

The effect of scale variations on the NLO+NNLL result is presented in (c) 13 and (d) 100 TeV for the same \(M_{W_R}\). The effect of the resummation on the scale variation is manifest in the reduction of the associated uncertainty. For the 13 TeV case, uncertainty is reduced to the sub per-cent level, while at 100 TeV the impact is comparable (but smaller) than the NLO dependence. This is because resummed contributions are less important away from threshold. Indeed, the observed reduction in scale uncertainty is consistent with what one expects from including higher order terms in the perturbative series.

We calculate the symmetric PDF uncertainties from the NNPDF member sets following the recommended procedure of Ref. \cite{r84}. The 68% (1\(\sigma\)) uncertainty bands are represented by the shaded regions in Fig. 3. In the upper panel, only the NLO+NNLL uncertainty are shown; in the lower panel, both the NLO and NLO+NNLL uncertainties are shown. At 13 TeV, for \(M_{W_R} = 4 \,(4.5)\) TeV, the NLO+NNLL uncertainty is approximately \(\pm 80 \,(240)\%\). At 100 TeV, the uncertainties breach 100% for \(M_{W_R}\) between 20 and 30 TeV.

The larger uncertainties in the threshold calculation compared to the NLO result is due in part to the less data used to constrain the threshold-improved PDFs \cite{r69, r70}. This follows from the limited threshold calculations available for processes that the enter into global fit PDFs, and demonstrates their need for accurate LHC predictions.

For representative \(M_{W_R}\), scale and PDF uncertainties are given in Tb. III.

IV. OBSERVABILITY OF BOOSTED \(N\) AT HADRON COLLIDERS

In this section we study the observability at hadron colliders of \(W_R\) and \(N\) in the LRSM for \(m_N/M_{W_R} \lesssim 0.1\). We start with production- and decay-level kinematics of \(N\) at LO. After constructing several observables with strong background-discriminating power, we perform a full parton shower (PS)/detector-level signal-to-background analysis.

For signal event generation, we modify the Manifest LRSM FeynRules (FR) model file v1.1.6.mix by Ref. \cite{r72} (see App. \textit{I}) and use FR v2.3.10 \cite{r101, r102} to generate Universal File Object (UFO) inputs \cite{r103}. LO events are simulated using MG5 \cite{r92}. Rates are scaled by the NLO+NNLL \(K\)-factors as defined in Eq. \(47\). Application of \(K\)-factors is justified in the threshold regime as the dominant contribution, i.e., soft-radiation, largely leave kinematics unchanged. Events are showered using PYTHIA 8.212 \cite{r104} and jets are clustered with FastJet v3.2.0 \cite{r105, r106} using the Cambridge/Aachen (C/A) algorithm \cite{r107, r108} with a separation parameter of \(R = 1.0\). SM background processes are simulated at LO+PS accuracy using the MG5, and scaled by an appropriate NLO \(K\)-factor calculated via the MG5-aMC@NLO framework. Due to extreme phase space cuts, event generation at NLO+PS accuracy is impractical.
To investigate the kinematics of boosted $N$ from $W_R$ decays, we simulate at 13 TeV

$$q_1 \bar{q}_2 \rightarrow W_R \rightarrow e_1 N \rightarrow e_1 e_2 q'_1 \bar{q}'_2,$$

(78)

where the two electrons possess any electric charge combination, for the representative $(M_{W_R}, m_N)$ listed in Tb. II. We focus on final-state electrons, which is the most problematic channel for ATLAS and CMS [55, 56], but our study is also applicable to the $e\mu$ and $\mu\mu$ final states. The largest change in those channels follows from the better
Inclusion of the $N \rightarrow ℓtb$ final state is similarly straightforward. To model detector response while keeping generator-level particle identification at LO, we smear final-state partons as done in Ref. [109], which adopts the expected ATLAS detector performance parametrization [110]. Eq. (78) is free of kinematic poles and no generator-level cuts are applied.

In Fig. 6 we show the normalized differential distributions with respect to the (a) transverse momentum ($p_T$) and (b) pseudorapidity ($η$) of the charged lepton in the $W^±R → Ne^±$ decay, denoted by $ℓWR$. In the $p_T^{ℓWR}$ distribution, the Jacobian peak near $p_T \sim M_{WR}/2$ is unambiguous and is largely independent of such small $m_N$. The $η^{ℓWR}$ distribution reveals that $ℓWR$ are very central, with most electrons contained within $|η| < 1.0$ and negligibly few with $|η| ≥ 2.0$.

Multi-TeV bounds on $M_{WR}$ (see Sec. II) nearly guarantee that $p_T^{ℓWR}$ is very large and $|η|$ small. Consequently, Eq. (78) efficiently passes inclusive high-$p_T$ single-electron triggers, such as those used in Ref. [62].

As $pp → N e^±$ is a 2 → 2 system, the heavy neutrino’s $p_T$ and rapidity ($y$) distributions are identical to Fig. 6 up to mass corrections. Hence, the decay products of the $N$ with high-$p_T$ are largely collimated due to its relative lightness.

FIG. 6. Normalized (a) transverse momentum ($p_T$) and (b) pseudorapidity ($η$) distributions of the charged lepton from $pp → W_R → Nℓ$ for representative $M_{WR}$ and $m_N$ at 13 TeV.

FIG. 7. Normalized distributions with respect to (a) $ΔR_{ℓq}$ and (b) $ΔR_{qq'}$ of $N$’s decay products for the same configuration as Fig. 6.
FIG. 8. Normalized distributions with respect to (a) \( p_T^{j_N} \), (b) \( y^{j_N} \), (c) \( m_{j_N} \), and (d) MET for same configuration as Fig. 6 but requiring exactly one electron and \( j_N \) candidate.

For the \( N \to \ell_N q q' \) final state in Eq. (78), we show in Fig. 7 the normalized separation distributions between (a) the charged lepton \( \ell_N \) and its closest quark (\( \Delta R_{qX}^{\text{min}} \)), as well as (b) the two quarks themselves (\( \Delta R_{qq'} \)). In both cases, the separation peaks at \( \Delta R \sim 0.2 \) (0.4) for \( r_N = m_N/M_{W_R} = 0.05 \) (0.1), and follows from the scaling relationship

\[
\Delta R_{qX} \sim 2 p_X^T / p_T^{j_N} \sim 4 m_N / M_{W_R},
\]

where \( p_X^T \) is the perpendicular momentum of \( X = \ell_N, q' \) relative to its parent \( N \). Hence, for much of the phase space, these electrons fail particle identification criteria at 13 TeV [62]:

\[
p_T^{j_N} > 35 \text{ GeV}, \quad \Delta R_{\ell X} > 0.3, \quad |\eta^\ell| < 2.4,
\]

and leads to the breakdown of current ATLAS and CMS \( W_R - N \) search strategies [55]. Smaller \( r_N = m_N^2 / M_{W_R}^2 \), hadronization, and the presence of a \( tb \) pairs exacerbate this issue.

---

1 The separation between particle pair \((a, b)\) is defined as \( \Delta R_{ab} \equiv \sqrt{(y_a - y_b)^2 + (\phi_a - \phi_b)^2} \) for rapidity \( y \) (or pseudorapidity \( \eta \)) and azimuthal angle \( \phi \).
For such signal regions, we consider an alternate search strategy: model the $N$ decay products as a single object, which we call a *neutrino jet* ($j_N$), and investigate instead the $2 \to 2$ process:

$$pp \to W_R \to e^\pm j_N.$$  

(81)

The simplified signal topology alleviates the failing identification criteria and retains the high signal-to-noise properties of the same-sign dilepton channel. To build a qualitative picture of the new signal definition, we preliminarily define $j_N$ at the present FO parton-level via C/A clustering with $\Delta R = 1$. We cluster all final-state partons except any electron candidate satisfying Eq. (80). $j_N$ is identified as the highest $p_T$ C/A jet.

In Fig. 8, we show the normalized distributions for $j_N$ with respect to (a) $p_T$, (b) $y$, (c) invariant mass ($m_{j_N}$), and (d) missing transverse momentum (MET) for events with exactly one electron and $j_N$ candidate. As anticipated, we observe strong similarities to the $\ell W_R$ distributions and unambiguous Breit-Wigner resonances at the appropriate values in the $m_{j_N}$ distribution. This indicates that $j_N$ is a good description of $N$ and that the signal definition of Eq. (81) can be interpreted as Eq. (78) when $(m_N/M_{W_R}) < 0.1$.

A cost of this new signal definition is the loss of the unambiguous *smoking-gun* collider signature of two same-sign leptons and jets [37], which is intrinsically background-free up to detector effects as it violates $L$ by two units. However, inherited from the original definition is the fact that, up to detector and hadronization effects, the process has no MET as no light neutrinos exist in the final state. Requiring again exactly one electron candidate, we show in (d) the normalized MET distribution. Due to smearing, we find moderate MET out to 10s of GeV and largely independent of $m_N$. We observe that the peak MET shifts to larger values for larger $M_{W_R}$ and is due to the increased likelihood of more energy being mis-reconstructed for more energetic objects [110]. Present ATLAS detector capabilities [111] permit MET cuts as tight as

$$\text{MET} < 35 \text{ GeV}.$$  

(82)

In a realistic scenario (see Sec. [IVB]), a more conservative cut is required due to pile up, etc.

In Fig. 9(a) the $W_R$ resonances built from the $\ell_1 - j_N$ invariant mass are clearly seen for our representative masses, up to broadening due to mis-reconstruction of $N$ and detector smearing. In (b), we show the polarization of $\ell_1$ in the $\ell_1 - j_N$ system’s rest frame. We observe clearly the RH chiral structure of the $N\ell W_R$ vertex for $\sqrt{r_N} = m_N/M_{W_R} = 0.01$. At larger $r_N$, however, this becomes obfuscated due to the importance of opposite helicity states, which scale like $r_N$, and lead to spin-decorrelation.

Altogether, this demonstrates the viability of the new search procedure.

Aside from the application of micro-jets and substructure techniques, it may be possible to verify the Majorana nature of heavier $N$ via its decays to top quarks. For Dirac $N$, the off-shell $W^*_R$ to which it decays can only carry the same electric charge as the charged lepton produced from the decay of the primary, on-shell $W_R$, i.e., $\ell W_R$. Decays of the $W^*_R$ to a top quark that subsequently decays leptonically can lead to final-state muons with the *same* sign electric
As the outgoing muon momentum scales like $q_{\mu}$, Hence, additional electrons fail to pass the criteria. Fake events correspond to regions of phase space where one electron candidate is identified according to Eq. (80) but $p_T^\ell > 35$ GeV + 2nd $e^\pm$ veto 

TABLE IV. Cross sections \([ab]\) of SM background for $pp \to e^\pm j_N$ after decays and successive cuts.

| Cut \(\sigma^{LO}\) [ab] | $W_j$ | $WZ$ | $t\bar{t}$ | $t\bar{t}j$ | $tbj$ | $eej$ | $WW_j$ |
|--------------------------|-------|------|---------|---------|-----|-----|------|
| $p_T^{j,b} > 30$ GeV, $|\eta^{j,b}| < 4.5$ | 218 | 218 | 218 | 218 | 218 | 218 | 218 |
| $+\Delta R_{j,b} > 0.4, \Delta R_{LX} > 0.3$ | 218 | 218 | 218 | 218 | 218 | 218 | 218 |
| No Decay | $10^6$ | $10^5$ | $10^6$ | $10^6$ | $10^6$ | $10^6$ | $10^6$ |
| +Decay+ | $p_T^{\ell, max} > 1$ TeV | 218 | 2.61 | 0.201 | 0.660 | 0.062 | 184 | 0.637 |
| $E_T < 50$ GeV | +Smearing+$|\eta^j| < 2.0$ | 218 | -- | -- | -- | 57 | -- |
| $p_T^\ell > 35$ GeV | 218 | -- | -- | -- | -- | -- | -- |
| $E_T < 35$ GeV | $R^{SM, R} = 1.3$ | 111 | -- | -- | -- | 25 | -- |
| $K_{SM, R} = 1.3$ | -- | -- | -- | -- | 33 | -- | -- |

charge as $\ell_{WR}$. That is, for a fixed primary $W_R$ electric charge, one has

$$q \bar{q} \to W_R^\pm \to \ell_{WR}^\pm N, \quad \text{with} \quad N \to \ell_N^- (t \to W^+_R b) b \to \ell_N^- b \bar{b} \mu^\pm \nu_{\mu}.$$  

Hence, $j_N$ containing top quarks can be identified by their larger complexity, namely the presence of two $b$-subjets. As the outgoing muon momentum scales like $p_T^\ell \sim \gamma_t m_t (1 + M^2_W/m_t^2)/4 \sim \gamma_t 50$ GeV, where $\gamma_t \sim m_t/p_T^N \sim m_t/M_W$ is the top quark’s Lorentz boost to the lab frame, it should be identifiable. For a Majorana $N$, the off-shell $W_R$ can carry either electric charge. Thus, observation of such muons with opposite electric charge of the easily identifiable $\ell_{WR}$ is evidence of $L$-violating transitions. Further discussion of this topic is beyond the scope of this study.

We briefly note that the use of neutrino jets is also widely applicable to other situations: In the MLRSM, high-mass $Z_R$ and $H_{FCNH}$ decays to boosted $N$ pairs could give rise to two back-to-back $j_N$. If $N$ couples non-negligibly to EW bosons, then $j_N$ may also feature substructure topologies. In other models, such as the Inverse Seesaw, rare decays of $W/Z/h$ bosons to GeV-scale pseudo-Dirac neutrinos, as well as other processes, could also result in $j_N$.

1. Estimation of Leading Standard Model Backgrounds

Before simulating our full detector-level analysis, we are in position to estimate the leading SM backgrounds. The simple lepton+jet topology of Eq. (81) suffers from large SM backgrounds. We sort the leading channels into three categories: (a) weak bosons, (b) top quarks, and (c) fake rates from electron misidentification:

Weak boson : $W^\pm j$ ($\to e^\pm j_X$), $W^\pm Z (\to e^\pm j_X)$ 

Top quark : $t\bar{t} + nj$ (semi-leptonic), $tbj (\to \ell^\pm + nb + mj_X)$

Fake rates : $e^+ e^- j$, $W^+ W^- j (\to e^+ e^- j_X)$

Fake events correspond to regions of phase space where one electron candidate is identified according to Eq. (80) but additional electrons fail to pass the criteria.

At the generator level and assuming the following (nominal) regulating cuts

$$p_T^{j,b} > 30 \text{ GeV}, \quad \Delta R_{j,b} > 0.4, \quad \Delta R_{LX} > 0.3, \quad |\eta^{j,b}| < 4.5,$$

the DY+1j channels at LO, i.e., $W_j$ and $eej$, are found to dominate with cross sections reaching $\sigma^{SM} \sim 0.3 - 2$ nb; see row 1 of Tb. [IV]. The signal/noise ratio roughly translates to $S/N \sim 10^{-6} - 10^{-5}$. Background rates are dramatically reduced after decaying the top quark and EW bosons, and requiring that the $p_T$ of the leading charged lepton and process MET satisfy at the generator level

$$p_T^\ell \text{ Generator-level} > 1 \text{ TeV} \quad \text{and} \quad \text{MET}^\text{Generator-level} < 50 \text{ GeV}.$$  

The $Wj$ and $eej$ channels remain dominant but now only reach $\sigma^{SM} \sim 200$ ab; see row 2 of Tb. [IV]. The top background is particularly neutralized owing to the cascade nature of their decays, which require TeV-scale charged leptons to be accompanied by TeV-scale light neutrinos from a multi-TeV top quark parent. Subsequently, the top quark and diboson backgrounds can be neglected.
TABLE V. $pp \rightarrow e^{\pm} j_{\text{fat}}$ rates [fb] after successive cuts and QCD normalization, as well as acceptance rate and statistical significance after all cuts for representative $(M_{W_R}, m_N)$ at $\sqrt{s} = 13$, 100 TeV.

| Cut | $\langle M_{W_R}, m_N \rangle$ [TeV, GeV] |
|-----|-----------------------------------|
|     | $(3, 30)$ | $(3, 150)$ | $(3, 300)$ | $(4, 400)$ | $(5, 500)$ |
| Fiducial+Kinematics + Detector+K-Factor [Eq. (90)] | 6.87 | 6.76 | 6.39 | 0.69 | 0.06 |
| MET [Eq. (91)] | 4.30 (63%) | 4.22 (62%) | 4.02 (63%) | 0.40 (58%) | 0.03 (50%) |
| $m_{\ell j_{\text{fat}}} [\text{Eq. (92)}]$ | 3.64 (85%) | 3.59 (85%) | 3.41 (85%) | 0.30 (75%) | 0.02 (67%) |
| $A = \sigma^{\text{eej}}_{\text{eej+Kin}}/\sigma^{\text{eej+Fat}}$ | 53% | 53% | 53% | 43% | 33% |
| $\sqrt{S/N}$ [10 fb$^{-1}$] | 5.9 | 5.9 | 5.7 | 1.7 | 0.4 |
| $\sqrt{S/N}$ [100 fb$^{-1}$] | 19 | 19 | 18 | 5.4 | 5.7 [2 ab$^{-1}$] |

$M_{W_R}$ and appropriate signal $K$-factor, the signal/noise ratio exceeds $S/N \gtrsim 10 – 100$.

B. Detector-Level Signal Analysis and Neutrino Jet Definition

Using a custom detector simulation, we model the effects of detector resolution and efficiency based closely on the ATLAS Kraków-parameterization $^{112}$. The parametrization provides a conservative estimate of the ATLAS detector performance for the phase-II high-luminosity LHC. We model pile-up (with $\mu = 80$) and $\Sigma E_T$-dependent resolutions for jets and MET. We define an electron to be isolated if the hadronic energy deposit within a cone of size $\Delta R < 10%$ of the lepton candidate’s $p_T$. For benchmark points we use the $(M_{W_R}, m_N)$ listed in Tb. III i.e., $m_N/m_{W_R} \lesssim 0.1$ at $\sqrt{s} = 13$ and 100 TeV. We summarize our analysis in Tb. V.

As described in Sec. IV A the angular separation between the charged lepton and the $W_R$ decay products in the chain $N \rightarrow \ell^\pm W^\mp \rightarrow \ell^\pm q\bar{q}'$, depends on the $W_R - N$ mass hierarchy. A significant amount of radiation from the $W_R$ decay enters the isolation cone of $\ell$ and can negatively affect the lepton’s identification. While so-called mini-isolation requirements $^{113}$ can be applied to recover the unidentified leptons, we adopt a more conservative approach and include the lepton’s momentum as part of a fat jet ($j_{\text{fat}}$), recombinated with the C/A algorithm and a cone size of $\Delta R = 1.0$, i.e., $j_N$. Hence, we focus on the inclusive process

$$pp \rightarrow W_R \rightarrow e^\pm N \rightarrow e^\pm j_{\text{fat}}.$$  \hfill (89)

We require the electron and $j_{\text{fat}}$ to further satisfy

$$p_T^e > 1 \text{ TeV}, \quad p_T^{j_{\text{fat}}} > 1 \text{ TeV}, \quad |\eta^e| < 2.5, \quad |\eta^j| < 2.5.$$ \hfill (90)

After kinematic and fiducial cuts, we see in row 5 (14) of Tb. V that the 13 (100) TeV rate for our representative $(M_{W_R}, m_N)$ spans 60 ab – 7 fb (0.2 – 1 pb). Including the detector response shifts the signal MET distribution to
In rows 6 and 15 of Tb. V, we find that about 50−Eq. (82) to larger values than estimated in Sec. IV A. In order to not lose a majority of the events, we loosen the MET cut of

\[ \text{MET} < 100 \text{ GeV}. \]  

In rows 6 and 15 of Tb. V we find that about 50 − 60% of events survive the MET requirement, with heavier (lighter) \( W_R \) having a lower (higher) survival likelihood. This behavior is due to the increase in momentum mis-measurement at larger \( p_T \) scales, which necessarily occurs with heavier \( M_{W_R} \), and is visible in the MET distribution of Fig. 8(d).

Similarly, higher collider energies lead to additional secondary radiation and larger MET.

In Fig. 10, we show the invariant mass distributions at LO+PS for the reconstructed heavy neutrino and (b) \( W_R \) including detector effects at LO+PS, as detailed in Sec. IV B.

We find that approximately 33−50% of events pass our selection criteria.

To further reduce the SM background, we apply the following cut around \( m_{\ell j_{\text{Fat}}} \):

\[ |m_{\ell j_{\text{Fat}}} - M_{W_R}| < 200 \text{ GeV}. \]  

The largeness of the mass window is motivated by the size of \( W_R \)'s total width \( \Gamma_{W_R} \). In row 7 (16) of Tb. V we see that roughly 60 − 85% (65 − 80%) of events at 13 (100) TeV rate pass this cut, again with heavier (lighter) \( W_R \) having a lower (higher) survival likelihood. The behavior here can be understood by comparing the 200 GeV mass window to \( \Gamma_{W_R} \) in Tb. III. For heavier (lighter) \( W_R \), we see that the mass window is about 1.4 (2.4) × \( \Gamma_{W_R} \), and hence encapsulating fewer (more) \( W_R \). As in the parton-level analysis, we find that the residual SM background is negligible.

For 13 (100 TeV), we calculate in row 8 (17) the acceptance rate, defined as the ratio of rows 7 and 5 (16 and 14):

\[ A \equiv \frac{\sigma^{\text{All Cuts}}}{\sigma^{\text{Fiducial+Kinematics+Detector Response}}}. \]  

We find that approximately 33 − 50% events pass our selection criteria.

Using the Gaussian estimator,

\[ \sigma = \frac{S}{\sqrt{S + B}} \approx \sqrt{S}, \quad \text{for} \quad S(B) = \mathcal{L} \times \sigma^{\text{All Cuts}} \text{ (SM background)} , \]  

we can determine the statistical significance of the signal process \( (S) \) over the SM backgrounds \( (B) \) after an integrated luminosity of \( \mathcal{L} \). At 13 TeV, we find a > 5\( \sigma \) statistical observation (discovery) for \( M_{W_R} = 3 \) (4) (5) TeV, independent of \( m_N \), after \( \mathcal{L} = 10 \) (100) (2000) fb\(^{-1}\). At 100 TeV and \( \mathcal{L} = 10 \) fb\(^{-1}\), all benchmark points are in excess of 20\( \sigma \). This is summarized in rows 9, 10, and 18 of Tb. V and Figs. III(a) and (b). We extrapolate the discovery potential for higher \( M_{W_R} \) by keeping fixed the efficiency, \( \varepsilon \equiv \sigma^{\text{Fid.+Kin.}} / \sigma^{\text{Total}} \), and acceptance for \( (M_{W_R}, m_N) = (5 \text{ TeV, } 500 \text{ GeV}) \), in which case \( \varepsilon \approx 0.64, \ A \approx 0.33 \). As seen in Fig. III(b) a 5\( \sigma \) discovery can be obtained for \( W_R \) masses up to \( M_{W_R} = 15 \) (30) with approximately 100 fb\(^{-1}\) (10 ab\(^{-1}\)).
Finally, we discuss briefly the 13 and 100 TeV potential to exclude previously unconstrained regions of the $(M_{W_R}, m_N)$ parameter space. We use Poisson counting to deduce the required luminosity $L_{95}$ for a 95% Confidence Level (CL) exclusion: For a SM background of $B \approx 0$ events, we solve for the largest number of signal events $S$ such that the expected probability to observe $B$ events is at most $5\% (= 1 - CL)$, i.e., find $S$ such that we satisfy:

$$\Pr (n_{\text{observed}} = B | n_{\text{expected}} = S + B) = \frac{(S + B)^B}{B!} e^{-(S+B)} \leq 1 - CL = 0.05.$$  

For $B \approx 0$, this yields $S = 3$. Given an efficiency $\epsilon$ and acceptance $A$, $L_{95}$ can then be determined by the relationship

$$L_{95} = \frac{S}{\epsilon \cdot A \cdot \sigma^{NLO+NNLL} \times BR \times BR}.$$  

We then state that a $(m_N/M_{W_R})$ mass-hypothesis is excluded at 95% CL if

$$N = L_{95} \times \sigma^{NLO+NNLL} \times BR \times BR \cdot \epsilon \cdot A \geq S = 3.$$  

For $(m_N/M_{W_R}) \leq 0.1$, we show in Fig. 11(c) that $M_{W_R} < 3$ (4) [5] TeV can be excluded at 95% CL with $L = 0.9 (10) \ [50 \text{ fb}^{-1}]$ of 13 TeV data. Also plotted are the ATLAS experiment’s 8 TeV 95% CL [55] and KamLAND-Zen 90% CL [31, 77] complimentary exclusion limits. We find that regions of the $(M_{W_R}, m_N)$ parameter space...
unconstrained by ATLAS and CMS are indeed covered by the present, complimentary analysis. The open region between this analysis and ATLAS is an artifact of our choice to limit our study to $(m_N/M_{W_R}) \leq 0.1$; application of the neutrino jet analysis to larger mass ratios will close the region. We note that the ability to exclude $M_{W_R} < 3$ TeV at $\sqrt{s} = 13$ TeV with approximately 1/20 of the 8 TeV data is consistent with the luminosity increase for DY-type processes [114]. In Fig. 11(d), we show the analogous 100 TeV exclusion potential: with $\mathcal{L} = 100$ fb$^{-1}$ (2.5 ab$^{-1}$), we find that $M_{W_R} < 22$ (33) TeV and $(m_N/M_{W_R}) \leq 0.1$ can be excluded at 95% CL.

V. SUMMARY AND CONCLUSION

The origin of tiny, nonzero neutrino masses remains an open question in particle physics. In this study, we re-examine the discovery potential of a $W_R$ gauge boson decaying to a heavy Majorana neutrino $N$ in the MLRSM. We focus on the case when $N$ is hierarchically lighter than $W_R$, i.e., $m_N/M_{W_R} \lesssim 0.1$. In this limit, $W_R \rightarrow N$ decays produce highly boosted $N$ that then decay to collimated final states. Subsequently, the canonical collider definition

$$p p \rightarrow W_R \rightarrow e^\pm N (\rightarrow e^\pm jj)$$

breaks down due to failing isolation criteria of the final-state charged leptons. For such a regime, we consider an alternative collider definition,

$$p p \rightarrow W_R \rightarrow e^\pm N \rightarrow e^\pm j_N,$$

where $j_N$ is a color-singlet neutrino jet and consists of the collimated $N$ decay products. Furthermore, we consider resummed QCD corrections that are important for high-mass DY processes. We calculate, for the first time, inclusive $pp \rightarrow W_R$ production at NLO+NNLL matched to threshold-improved PDFs. This captures dominant contributions beyond NLO, and are arguably the most precise predictions available for high-mass $W_R$ at 13 and 100 TeV. We summarize our findings:

1. We introduce the concept of neutrino jets, which has widespread applicability to other processes and models; see Sec. [V.A]. With our new collider signal definition, a $5 - 6 \sigma$ discovery is achievable at 13 TeV with 10 (100) [2000] fb$^{-1}$ for $M_{W_R} = 3$ (4) [5] TeV and $(m_N/M_{W_R}) < 0.1$. At 100 TeV, a $5 \sigma$ discovery can be obtained for $W_R$ masses up to $M_{W_R} = 15$ (30) TeV with approximately 100 fb$^{-1}$ (10 ab$^{-1}$). Conversely, with 0.9 (10) [150] fb$^{-1}$ of 13 TeV data, $M_{W_R} < 3$ (4) [5] TeV can be excluded at 95% CL; with 100 fb$^{-1}$ (2.5 ab$^{-1}$) of 100 TeV data, $M_{W_R} < 22$ (33) TeV can be excluded. See Sec. [V.B]

2. At 13 TeV, the NLO+NNLL contributions increase the Born (NLO)-level predictions by 40 − 140 (4 − 70)% for $M_{W_R} = 4 - 5$ TeV, well beyond the NLO scale uncertainty. At 100 TeV, threshold effects become important for $M_{W_R} \gtrsim 30$ TeV, where resummation increases the Born (NLO) prediction by $\gtrsim 40 (10)$%. Away from threshold, we find that the resummed result converges to the NLO rate. See Sec. [III]

3. The residual scale dependence at NLO+NNLL for $M_{W_R} = 1 - 5$ (1 - 30) TeV at 13 TeV is maximally sub-percent, and $\pm 4\%$ at 100 TeV. The PDF uncertainty at NLO+NNLL exceeds 100% in the threshold regions. Away from threshold, the PDF uncertainty is comparable to the NLO PDF uncertainty. See Sec. [III.D]

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Appendix A: Threshold Resummation for Inclusive $pp \rightarrow W'$ Production

Here, we review the details of threshold resummation for inclusive production of $W'$ bosons with arbitrary chiral couplings in $pp$ collisions. Often labeled as soft-gluon or large-$x$ resummation [93, 95], the calculation should not be confused with small-$k_T$ (or recoil or Collins-Soper-Sterman) resummation [115–117], nor joint recoil-threshold resummation [118–120]. Many concise texts on the topic exist, e.g., Refs. [96, 99, 121–123]. We largely follow the notation and spirit of Refs. [99, 98] and implement the numerical procedures of Refs. [96, 94], but make explicit mass and coupling factors that are often omitted for simplicity.
1. Threshold Resummation for $W'$ Bosons with Arbitrary Chiral Couplings

The emission of soft gluons from initial-state partons participating in hard, hadronic collisions can spawn large, with respect to the expansion parameter $\alpha_s$, logarithmic enhancements in scattering cross sections. In certain kinematic configurations, these logarithms can become large enough to render a perturbative expansion unreliable, requiring that the divergent series be summed to all orders. In particular, soft logarithms near the partonic threshold take the form

$$\alpha_s \log \left( \frac{\hat{s} - Q^2}{\hat{s}} \right) = \alpha_s \log(1 - z), \quad z = \frac{Q^2}{\hat{s}}, \quad (A1)$$

where $Q \sim \sqrt{\hat{s}} \gg \Lambda_{\text{QCD}}$ is the scale of the hard scattering process, $\sqrt{\hat{s}}$ is the partonic scattering scale, and the dimensionless variable $z$ quantifies the nearness of the partonic scale to the hard scale. A schematic distinction of the hard, partonic, and hadronic scattering (beam) scale $\sqrt{s}$ is illustrated in Fig. 2. The purpose of threshold resummation is to perform a summation of such terms when $z \to 1$ while accounting for the hierarchy of scales via renormalization group evolution (RGE). We now briefly summarize the procedure directly in perturbative QCD (pQCD).

For a generic color-singlet boson $B$ (scalar or vector) produced in hadron collisions, the total inclusive cross section is given by the usual Collinear Factorization Theorem

$$\sigma^{\text{FO}}(h_1 h_2 \to B + X) = \sum_{a,b=q,q',g} \int_{\tau_0}^{1} d\tau L_{ab}(\tau, \mu_f) \times \hat{\sigma}^{\text{FO}}(ab \to B), \quad \tau_0 \equiv \frac{M_B^2}{\hat{s}}, \quad (A2)$$

where the luminosity $L$ of parton pair $ab$, with $a,b \in \{q,q',g\}$, at the LHC ($h_1 = h_2 = p$) is given in terms of the PDFs $f_{a/p}$ and $f_{b/p}$ jointly evolved to a factorization scale $\mu_f$:

$$L_{ab}(\tau, \mu_f) = \frac{1}{1 + \delta_{ab}} \int_{\tau}^{1} d\xi_1 \frac{d\xi_1}{\xi_1} \left[ f_{a/p}(\xi_1, \mu_f) f_{b/p}(\xi_2, \mu_f) + f_{a/p}(\xi_2, \mu_f) f_{b/p}(\xi_1, \mu_f) \right], \quad (A3)$$

$$\xi_2 \equiv \frac{\tau}{\xi_1}, \quad (A4)$$

and $\hat{\sigma}^{\text{FO}}$ is the FO partonic cross section for the process

$$a \ b \to B \quad \text{with} \quad Q = M_B. \quad (A5)$$

Following the notation and methodologies of Refs. 96, 98, we account for the arbitrary emission of soft radiation by using a generalization of Eq. (A2):

$$\sigma^{\text{FO}}(pp \to B + X) = \sum_{a,b=q,q',g} \int_{\tau_0}^{1} d\tau \int_{0}^{1} dz \, \delta \left( z - \frac{\tau_0}{\tau} \right) \times L_{ab}(\tau) \times \hat{\sigma}^{\text{FO}}(ab \to B). \quad (A6)$$

Written this way, we identify the FO partonic cross section in the soft radiation limit as

$$\hat{\sigma}^{\text{FO}}(ab \to B) \equiv \hat{\sigma}^{\text{FO}}_{ab} = \sigma_0 \times z \times \Delta^{\text{FO}}_{ab}(z), \quad \Delta^{\text{FO}}_{ab}(z) = \sum_{k=0}^{\infty} \left( \frac{\alpha_s}{\pi} \right)^k \Delta^{(k)}_{ab}(z). \quad (A7)$$

The soft threshold coefficient $\Delta_{ab}(z)$, which encapsulates the factorizable soft emissions, is often denoted as $C_{ab}$ and $G_{ab}$ in literature. For a $W'$ gauge boson with arbitrary chiral couplings $g_L$ and $g_R$ to quarks and mass $M_B = M_{W'}$, the constant term is

$$\sigma_0 = |V_{\text{CKM}}^{C}\left( g_L^2 + g_R^2  \right) \pi |^{2} \frac{4 N_c M_{W'}^2}{4 N_c M_{W'}^2}. \quad (A8)$$

We suppress the indices on $\sigma_0$ as the trivial generalization introduces an unnecessary notational complication. The expression is related to the usual LO partonic expression by

$$\hat{\sigma}^{\text{LO}}(q \ q' \to W') = \sigma_0 \times M_{W'}^2 \times \delta (\hat{s} - M_{W'}^2) = \sigma_0 \times z \times \delta(1 - z). \quad (A9)$$

At LO, one may identify $\Delta$ with the above $\delta$-function, which is determined by kinematics alone. This is because the LO $2 \to 1$ process occurs identically at threshold. Beyond LO, the structure of soft logarithms in the perturbative
expansion of $\Delta(z)$ remains essentially kinematic in origin \footnote{122}. In terms of explicit scale dependence, one can also write Eq. \footnote{17} as
\[
\delta_{ab}^{\text{FO}} = z \times \Delta_{ab}^{\text{FO}}(\hat{s}, Q^2) \times \sigma_0(Q^2 = M_0^2).
\]  
(A10)

This suggestive form indeed implies, in the language of RGE, that $\Delta(\hat{s}, Q^2)$ is an evolution operator that runs the hard process at $Q^2 = M_0^2$, up to the partonic scale $\hat{s} = \tau_s$ \footnote{122}.

If working with pQCD, the threshold resummed cross section can be efficiently obtained after writing the hadronic cross section in so-called Mellin-(or $\tau$-)space. This is because such convolutions become products in Mellin space. Applying the Mellin transform, as defined in Eq. (40), to Eq. (A6) with respect to expansion of $\Delta(\hat{s}, Q^2)$ is given as an expansion in $S$ \footnote{122}.

Among other considerations, the resummation is usually performed in the large-\(N\) limit as the $N \to \infty$ limit corresponds to the $z \to 1$ (threshold) limit for partonic cross sections. In this limit, additional gluon emission is constrained to be soft, and is therefore exactly where one finds a perturbative expansion rendered unreliable by large logarithms. Specifically, the divergent contributions at leading power in $(1-z)$ are plus distributions of the form
\[
\Delta^{(j)}(z) \sim \alpha_s^j(Q^2) \left[ \frac{\log^m(1-z)}{1-z} \right]_+, \quad m \leq 2j - 1.
\]  
(A13)

In Mellin space and in the large-$N$ limit, such distributions are transformed to a series of the form
\[
\Delta^{(j)}_N(z) \sim \alpha_s^j(Q^2) \sum_{r=0}^{2j} b_r \log^r N,
\]  
(A14)

where $b_r$ is some $N$ independent coefficient. To all orders in $\alpha_s$, resummation captures a number of these divergent logarithms, producing a finite result that can supplement FO calculations. For the $k$th term in the expansion, resummation corresponds at leading log (LL) accuracy to gathering all logarithms with power of $r = 2k$; at next-to-leading log (NLL), all logs such that $2k \geq r \geq 2k - 2$; and generically at $\text{N}^{(k)}\text{LL}$, $2k \geq r \geq 2k - 2j$. Furthermore, this implies that in re-expanding $\text{N}^{(k)}\text{LL}$ in $\alpha_s$, one can identify the inclusive $\text{N}^{(j-1)}\text{LO}$ calculation in the limit where all radiation is soft. This necessitates a matching scheme when combining resummed and FO results beyond LO.

In the notation of Ref. \footnote{97}, the resummed coefficient $\Delta_{q\bar{q}}^{\text{Res}}$ for color-singlet $q\bar{q}$ pairs is
\[
\Delta_{q\bar{q}}^{\text{Res}}(N) = g_0(\alpha_s) \exp S(\lambda, \mu), \quad \text{with} \quad \lambda = \mu \log \frac{1}{N} \quad \text{and} \quad \mu = a \alpha_s(Q^2) \beta_0,
\]  
(A15)

where $a = 2 (1)$ for DY (DIS) accounts for the number of contributing initial-state hadrons, and the $\text{Sudakov factor} S$ is given as an expansion in $\mu$, while treating $\mu \ln N \sim \mathcal{O}(1)$:
\[
S(\lambda, \mu) = \sum_{m=0}^{\infty} \mu^{m-1} g_{m+1} = \frac{1}{\mu} g_1(\lambda) + g_2(\lambda) + \mu g_3(\lambda) + \mathcal{O}(\mu^2) \equiv \sum_{k=0}^{\infty} \alpha_s^k S_k.
\]  
(A16)
We note that it is possible to consistently re-expand $\mathcal{S}$ in terms of $\alpha_s$ and coefficients $\mathcal{S}_k$. The normalization function $g_0$ is similarly perturbative and is given by

$$g_0(\alpha_s) = \sum_{n=0}^{\infty} g_{0n} \alpha_s^n = g_{00}^{\text{LL}} + \alpha_s g_{01}^{\text{LL}} + \alpha_s^2 g_{02}^{\text{LL}} + \mathcal{O}(\alpha_s^3)$$  \hspace{1cm} (A17)

Expressions and normalizations for $g_m$, $g_{0n}$, and the QCD $\beta$-function coefficient $\beta_0$ are detailed in [97]. Acquiring a resummation of order $\text{N}_3^\text{LL}$ is achieved by including the matching functions in $g_0(\alpha_s)$ up to $\mathcal{O}(\alpha_s^2)$, i.e., all $g_{0n}$ up to $n = j$, and $g_m$ functions for $m$ up to $m = j + 1$. In this work, we resum soft radiation up to next-to-next-to-leading logarithmic (NNLL) accuracy. Thus our resummed soft function is

$$\Delta_{(q\bar{q})N}^{\text{NNLL}} = (g_{00} + g_{01}\alpha_s + g_2\alpha_s^2) \exp \left[ \frac{1}{\pi} g_1(\lambda, \pi) + g_2(\lambda, \pi) + \pi g_3(\lambda, \pi) \right],$$  \hspace{1cm} (A18)

and our resummed cross section in Mellin-space at NNLL

$$\sigma_{N}^{\text{NNLL}}(pp \to W') = \sigma_0 \mathcal{L}_{q\bar{q}}(N+1) \Delta_{(q\bar{q})/(N+1)}^{\text{NNLL}}.$$  \hspace{1cm} (A19)

2. Inverse Mellin Transformation via Minimal Prescription Procedure

Taking the inverse Mellin transformation of Eq. (A19), as defined in Eq. (11), gives the resummed production cross section in momentum space:

$$\sigma_{\text{Res}}(pp \to W' + X) = \frac{\sigma_0}{2\pi i} \int_{c-i\infty}^{c+i\infty} dN \, t_0^{-N} \times \mathcal{L}_N \times \Delta_{\text{Res}}.$$  \hspace{1cm} (A20)

Formally, the integration path, with $c \in \mathbb{R}$, is to the right of all singularities. In practice, this is impossible due to the QCD Landau pole at $N = N_L \equiv \exp[1/2\alpha_s \beta_0]$. The situation can be remedied by adhering to the Minimal Prescription (MP) procedure [96], which entails choosing $c = C_{\text{MP}}$ such that

$$2 < C_{\text{MP}} < N_L,$$  \hspace{1cm} (A21)

to avoid the pomeron (Landau) pole as small (large) $N$, and deforming the path toward $\text{Re}[N] < 0$. The path deformation is illustrated in Fig. 12(b). Subsequently, Eq. (A20) becomes

$$\sigma_{\text{Res}}(pp \to W' + X) = \frac{\sigma_0}{2\pi i} \int_{C_{\text{MP}}} dN \, t_0^{-N} \times \mathcal{L}_N \times \Delta_{\text{Res}}.$$  \hspace{1cm} (A22)

$$= \frac{\sigma_0}{\pi i} \int_{C_{\text{MP}}[\text{Im}[N] > 0]} dN \, t_0^{-N} \times \mathcal{L}_N \times \Delta_{\text{Res}}.$$  \hspace{1cm} (A23)

In the second line, a factor of 2 follows from the integrand being even with respect to $\text{Im}[N]$. This follows from the fact that the original cross section in momentum space Eq. (A6) is a real function.

Following Ref. [124], and the associated code ResHiggs, we choose the path,

$$N(t) = C_{\text{MP}} + (m_{\text{MP}} - i) \log(t), \quad t \in (0, 1),$$  \hspace{1cm} (A24)

where $C_{\text{MP}}$, $m_{\text{MP}} \in \mathbb{R}$ cannot be too large numerically without hitting machine precision limitations in $t_0^{-N} = \exp[-(N - 1) \log t_0]$. We have checked $5 < C_{\text{MP}} < 15$ and $m_{\text{MP}} = C_{\text{MP}}/10$ leaves the integral unchanged. Making the change of variable to $t$, one has

$$\sigma_{\text{Res}}(s) = \frac{\sigma_0}{\pi} \int_0^1 \frac{dt}{t} (i - m) \, t_0^{-N} \times \mathcal{L}_N \times \Delta_{\text{Res}}.$$  \hspace{1cm} (A25)

$$= \frac{\sigma_0}{\pi} \int_0^1 \frac{dt}{t} \text{Im} \left[ (i - m) \, e^{-(N-1) \log t_0} \times \mathcal{L}_N \times \Delta_{\text{Res}} \right].$$  \hspace{1cm} (A26)

In the last line, we use the fact that $\sigma_{\text{Res}}$ is a physical rate, i.e., positive-definite, implying that the integrand must be purely imaginary to cancel the $1/i$.

\[2\text{ I.e., } N^3\text{LL in the } ^{\ast} \text{ convention or } N^\prime \text{LL}‘ \text{ in the } ^{\ast} \prime \text{ convention, which are precisely defined in Ref. [99].} \]
3. Matching Resummed and Fixed Ordered Expressions

The resummation procedure is derived in the threshold limit to leading power in $(1 - z)$ and therefore neglects sub-leading power corrections, e.g., hard, wide-angle radiation. To make viable predictions at collider experiments, it is desirable to supplement resummed formulae with exact FO results beyond LO. This is achievable by subtracting from the resummed expression the $O(\alpha_s^k)$ radiation terms common to both calculations, to avoid double counting of soft radiation, and add back the full FO $O(\alpha_s^k)$ result, which describes both soft and hard radiation. A convenient way to isolate those common terms is to Taylor expand the resummed expression $\sigma_{\text{Res.}}$. About $\alpha_s = 0$. Up to NLO, this is given by the first two terms of the expansion

$$\sigma_{\text{Res.}} = \sum_{l=0}^{\infty} \frac{\alpha_s^l}{l!} \left[ \frac{\partial^l}{\partial \alpha_s^l} \sigma_{\text{Res.}} \right]_{\alpha_s = 0}$$

(A27)

As the Mellin- and inverse-Mellin operators commute with the $\partial_{\alpha_s}$ operator, the expansion holds in Mellin-space. Furthermore, as there is no explicit $\alpha_s$ dependence in $\sigma_0$ and $L_N$ in Eq. (A19), the expansion of $\sigma_{\text{Res.}}$ is simply proportional to the Taylor expansion of the exponentiated coefficient $\Delta_{\text{Res.}}^N$. In Mellin space, the $O(\alpha_s)$-subtracted resummed cross section for $pp \rightarrow W_R$ is then

$$\sigma_{\text{Res.}}_{\alpha_s-\text{Subtracted}} = \sigma_0 \times L_N \times \left[ \Delta_{\text{Res.}}^N - \Delta_{\text{Res.}}^N \right]_{\alpha_s=0} - \alpha_s \left[ \partial_{\alpha_s} \Delta_{\text{Res.}}^N \right]_{\alpha_s=0}$$

(A29)

For color-singlet $q\bar{q}$ initial states, explicit calculation shows

$$\Delta_{\text{Res.}}^{(q\bar{q})N} \bigg|_{\alpha_s=0} = g_{00} \quad \text{and} \quad \partial_{\alpha_s} \Delta_{\text{Res.}}^{(q\bar{q})N} \bigg|_{\alpha_s=0} = (S_1 + g_{01}),$$

(A30)

which correspond to terms in Eq. (A10) and can be found in [97]. To NLO+NNLL accuracy, the physical inclusive $W'$ production cross section in $pp$ collisions is at least

$$\sigma_{\text{NLO+NNLL}}(pp \rightarrow W' + X) = \sigma_{\text{NLO}} + \sigma_{\text{NNLL}} \bigg|_{\alpha_s-\text{Subtracted}},$$

(A31)
where the modified resummed term is obtained by inserting Eq. \(A29\) into Eq. \(A26\):

\[
\sigma^{\text{NNLL}}_{\alpha_s-\text{Subtracted}} = \frac{\sigma_0}{\pi} \int_0^1 \frac{dt}{t} \text{Im} \left[ (i - m) \tau_{0}^{-2(N-1)} \times \mathcal{L}_N \times \left( \Delta^{\text{NNLL}}_{(qq')N} - g_{00} - \alpha_s(S_1 + g_{01}) \right) \right].
\]  

(A32)

4. Parton Luminosities in Mellin Space

The resummation formalism we exploit requires parton luminosities \(L_{qq'}\) in Mellin space. This introduces a technical difficulty as modern PDF sets are typically only available numerically. It is possible, however, at a fixed factorization scale \(\mu_f\), to approximate luminosities \(L_{ab}(\tau)\) [and individual PDFs \(f_{a/p}(\xi)\)] using a basis of polynomials that can be Mellin-transformed analytically. Here we use a basis of Chebyshev polynomials of the first kind \(T_n(x)\), for which fast numerical algorithms exist for calculating the expansion coefficients, e.g. [125]. The implementation and optimization of the Chebyshev approximation procedure has been documented in Refs. [97, 99]. We briefly summarize for completeness how to obtain the approximated \(L_{(qq')N}\).

We write a general Chebyshev polynomial of degree \(n\) defined over the domain \(x \in (-1, 1)\) as

\[
T_n(x) = \sum_{m=0}^{n} T_{mn} x^n, \quad T_{mn} \in \mathbb{Z}.
\]

The function \(F(u)\) over the domain \(u \in (u_{\text{min}}, u_{\text{max}})\) can then be approximated by the first \(n_{\text{ch}}\) polynomials by the relationship

\[
F(u) \approx \frac{c_0}{2} + \sum_{k=0}^{n_{\text{ch}}} c_k T_k(Au + B),
\]

(A33)

with \(A\) and \(B\) given by

\[
A = \frac{2}{u_{\text{max}} - u_{\text{min}}} \quad \text{and} \quad B = -\frac{u_{\text{max}} + u_{\text{min}}}{u_{\text{max}} - u_{\text{min}}},
\]

(A34)

and the \(k\)th Chebyshev coefficient \(c_k\) by [125]

\[
c_k = \frac{2}{n_{\text{ch}} + 1} \sum_{j=0}^{n_{\text{ch}}} \tilde{F}_j \cos \left( \frac{k\pi(j + \frac{1}{2})}{n_{\text{ch}} + 1} \right).
\]

(A35)

The \(j\)th moment of \(F(u)\), i.e., \(\tilde{F}_j\), is defined as

\[
\tilde{F}_j = F(y_j), \quad \text{with} \quad y_j = \frac{1}{2} \left( u_{\text{max}} - u_{\text{min}} \right) \cos \left( \frac{\pi(j + \frac{1}{2})}{n_{\text{ch}} + 1} \right) + \frac{u_{\text{max}} + u_{\text{min}}}{2}.
\]

(A36)

Such efficient algorithms allow us in principle to immediately obtain the luminosity \(L_{qq'}(\tau, \mu_f)\) in Mellin space by transforming Eq. \(A33\) directly. However, \(L(\tau)\) is generally poorly behaved across \(\tau \in (0, 1)\), particularly at the origin. This is resolvable by approximating a regularized version of the luminosity and set

\[
F(u) = \tau(u)L_{qq'}(\tau(u), \mu_f), \quad \text{with} \quad \tau(u) = e^u \quad \text{for} \quad u \in (\log \tau_0, 0).
\]

(A37)

As defined in Eq. \(A39\), \(\tau_0 = M_h^2/s\) is the threshold above which \(pp \rightarrow W'\) is kinematically allowed to proceed. After a wee bit of algebra, we obtain an expression for the Mellin transformed parton luminosities,

\[
L_{(qq')N} = \int_0^1 d\tau \, \tau^{N-1} \, L_{qq'}(\tau, \mu_f) = \int_0^1 d\tau(u) \, \tau^{N-2}(u) \, F(u) = \sum_{p=0}^{n_{\text{ch}}} \frac{\bar{c}_p}{(N-1)^{p+1}},
\]

(A38)

where we have defined

\[
\bar{c}_p = \frac{2p}{u_{\text{min}}} \sum_{j=p}^{n_{\text{ch}}} j! (j - p)! \tilde{c}_j, \quad \text{with} \quad \bar{c}_j = -\frac{c_0}{2} \delta_{j0} + \sum_{k=j}^{n_{\text{ch}}} c_k T_{kj}.
\]

(A39)
Once one calculates the initial coefficients $c_k$, it is straightforward to approximate the Mellin transform of $L_{qf}(\tau, \mu_f)$ by using Eqs. (A38)-(A39). However, for different $\mu_f$ choices, the function being approximated changes and therefore the coefficients $c_k$ need to be recomputed. This should be taken into account if one intends to use a dynamic factorization scale.

Appendix B: Modeling Manifest Left-Right Symmetric Model with FeynRules

The most generic scalar potential of the LRSM consists of 18 parameters: three mass scales $\mu_1, \ldots, \mu_3$, 14 dimensionless couplings $\lambda_1, \ldots, \lambda_4, \rho, \rho_\Phi$, and one CP violating phase $\delta_2$. It is given by \[21, 126\],

$$V(\Phi, \Delta_L, \Delta_R) = V_\mu + V_\Phi + V_{\Delta} + V_{\Phi\Delta} + V_{\Phi\Delta L, \Delta_R},$$

(B1)

where the scalar mass and self-coupling terms of the bidoublet $\Phi$ are, respectively,

$$V_\mu = -\mu_1^2 \text{Tr} \left[ \Phi^\dagger \Phi \right] - \mu_2^2 \text{Tr} \left[ \Phi^\dagger \Phi + \Phi \Phi^\dagger \right] - \mu_3^2 \text{Tr} \left[ \Delta_L \Delta_L + \Delta_R \Delta_R \right],$$

(B2)

$$V_\Phi = \lambda_1 \left[ \text{Tr} \left[ \Phi^\dagger \Phi \right] \right]^2 + \lambda_2 \left[ \text{Tr} \left[ \Phi^\dagger \Phi \right] \right]^2 + \lambda_3 \left[ \text{Tr} \left[ \Phi^\dagger \Phi \right] \right]^2$$

$$+ \lambda_4 \text{Tr} \left[ \Phi^\dagger \Phi \right] \text{Tr} \left[ \Phi^\dagger \Phi + \Phi \Phi^\dagger \Phi \right].$$

(B3)

The $\Delta_{L,R}$ self- and cross couplings are:

$$V_\Delta = \rho_1 \left[ \text{Tr} \left[ \Delta_L \Delta_L \right] \right]^2 + \rho_2 \left[ \text{Tr} \left[ \Delta_L \Delta_R \right] \right]^2 + \rho_3 \text{Tr} \left[ \Delta_L \Delta_L \right] \text{Tr} \left[ \Delta_R \Delta_R \right]$$

$$+ \rho_4 \text{Tr} \left[ \Delta_L \Delta_R \right] \text{Tr} \left[ \Delta_R \Delta_R \right] + \rho_5 \text{Tr} \left[ \Delta_L \Delta_R \right] \text{Tr} \left[ \Delta_L \Delta_R \right].$$

(B4)

The $\Phi - \Delta_L$ and $\Phi - \Delta_R$ couplings are

$$V_{\Phi \Delta} = \beta_1 \text{Tr} \left[ \Phi^\dagger \Delta_L \Phi \Delta_R + \Delta_L \Phi^\dagger \Delta_R \Phi \right] + \beta_2 \text{Tr} \left[ \Phi^\dagger \Delta_L \Phi^\dagger \Delta_R \Phi \right]$$

$$+ \beta_3 \text{Tr} \left[ \Phi^\dagger \Delta_L \Phi^\dagger \Delta_R \Phi^\dagger \Delta_R \Phi \right].$$

(B5)

After LR and EWSB, there exists 10 physical scalars: four neutral, CP even states $H_{0,\pm}^\pm$, including one at $m_{H_0^0} \approx 125$ GeV; two neutral CP odd states $A_{0,1}^0$; two states singly charged under $U(1)_E$ $H_{1,2}^\pm$; and two doubly charged states $\delta_{L,R}^\pm$. Subscripts do not indicate a mass ordering. The mass spectrum in the vev limit of Eq. (12) is given by \[21, 126\]:

$$m_{H_0^0}^2 \approx (125 \text{ GeV})^2 \approx 2k_+^2 \left( \lambda_1 + 4\frac{k_+^2 k_-^2}{k_+^2} (2\lambda_2 + \lambda_3) + 4\lambda_4 \frac{k_+^2 k_-^2}{k_+^2} \right).$$

$$M_{H_0^0}^2 = M_{H_1^\pm}^2 \approx \alpha_3 \frac{v_R^2}{2} \frac{k_+^2}{k_+^2}, \quad M_{H_0^0}^2 = M_{H_2^\pm}^2 \approx (\rho_3 - 2\rho_1) \frac{v_R^2}{2} k_+^2, \quad M_{H_0^0}^2 \approx 2\rho_1 v_R^2,$$

$$M_{H_1^\pm}^2 \approx (\rho_3 - 2\rho_1) \frac{v_R^2}{2} + \alpha_3 \frac{k_+^2}{2}, \quad M_{H_2^\pm}^2 \approx (\rho_3 - 2\rho_1) \frac{v_R^2}{2} + \alpha_3 \frac{k_+^2}{2},$$

$$M_{H_1^\pm}^2 \approx (\rho_3 - 2\rho_1) \frac{v_R^2}{2} + \alpha_3 \frac{k_+^2}{4}, \quad M_{H_2^\pm}^2 \approx 2\rho_1 v_R^2 + \alpha_3 \frac{k_+^2}{2},$$

(B7)

where $k_\pm$ is defined in Eq. (4).

With choice assumptions, the potential can be configured such that the theory is consistent with experimental limits and features new gauge states accessible by the LHC or VLHC. Accordingly, the Manifest LRSM FeynRules model of \[73\] can be set to simulate this region of the MLRSM parameter space. We now discuss this configuration.
1. Phenomenological Constraints on LRSM Scalar Potential

Explicit CP conservation and minimization conditions of the potential give rise to the so-called vev Seesaw relationship [120]:

\[ v_L = \frac{\beta_2 k_2^2 + \beta_1 k_1 k_2 + \beta_3 k_3^2}{(2\rho_1 - \rho_3)v_R}, \]

implying inherently small \( v_L \) for \( 2\rho_1 \neq \rho_3 \) and \( v_R^2 \gg k_1, k_2 \). Though it is natural for all \( \rho_i \) to be comparable in magnitude, it is contrived to expect a fine cancellation, particularly after radiative corrections. Consistent with \( \rho/T \)-parameter measurements [22, 23], we choose

\[ v_L = 0 \iff \beta_{1,2,3} = 0. \]  

(B9)

This may also be achievable if Eq. (B1) respects an approximate custodial symmetry.

Neglecting terms \( \mathcal{O}(k_2^2/v_R^2) \), the minimization conditions also imply [120]

\[ \frac{\lambda^2}{v_R^2} = \alpha_1 \frac{\alpha_3}{2} \left( \frac{k_2^2}{k_2^2} \right), \quad \frac{\lambda^2}{v_R^2} = \alpha_1 \frac{\alpha_3}{4} \left( \frac{k_1 k_2}{k_2^2} \right), \quad \frac{\mu_1^2}{v_R^2} = \rho_1. \]

(B10)

As argued, one expects on naturalness grounds

\[ \alpha_{2,3} \sim \mathcal{O}(\alpha_1) \quad \text{and} \quad \rho_{2,3} \sim \mathcal{O}(\rho_1). \]

(B11)

Dropping terms relatively suppressed by \( (k_2/k_\pm) \sim (m_6/m_i) \sim 10^{-2} \) [see Eq. (3)] gives

\[ \frac{\mu_1^2}{v_R^2} \approx \frac{\mu_2^2}{v_R^2} \approx \frac{\alpha_1}{2}, \quad \frac{\mu_3^2}{v_R^2} = \rho_1, \]

suggesting that LRSB is inherently at the mass scale of the scalar potential assuming

\[ \alpha_1 \sim \mathcal{O}(1) \quad \text{and} \quad \rho_1 \sim \mathcal{O}(1). \]

(B13)

In terms of \( M_{W_R} \) and \( g \), Eq. (B13) and positivity of squared masses for (physical) scalars imply several mass and coupling relationships:

\[ \frac{m_{H^0_{W_R}}^2}{M_{W_R}^2}, \frac{m_{A^0_{W_R}}^2}{M_{W_R}^2}, \frac{m_{H^\pm_{W_R}}^2}{M_{W_R}^2} \approx \frac{\alpha_3}{g^2} > 1, \quad \frac{m_{H^0_{W_R}}^2}{M_{W_R}^2} \approx \frac{4\rho_1}{g^2} > 1, \]

\[ \frac{m_{A^0_{W_R}}^2}{M_{W_R}^2}, \frac{m_{H^\pm_{W_R}}^2}{M_{W_R}^2} \approx \frac{(\rho_3 - 2\rho_1)}{g^2} > 0, \quad \frac{m_{A^0_{W_R}}^2}{M_{W_R}^2} \approx \frac{4\rho_2}{g^2} > 1. \]

(B14)

Imposing the strong requirement on Eq. (B14) to universally comply with bounds on FCNH, i.e., \( m_{FCNH} \) in Eq. (24), implies

\[ \rho_{1,2,4} > \frac{g^2}{4} \left( \frac{m_{FCNH}}{M_{W_R}} \right)^2, \quad \rho_3 > g^2 \left( \frac{m_{FCNH}}{M_{W_R}} \right)^2 + 2\rho_1 \approx 6\rho_1, \]

\[ \alpha_{1,2,3} > g^2 \left( \frac{m_{FCNH}}{M_{W_R}} \right)^2, \quad \mu_{1,2}^2 > (m_{FCNH})^2, \quad \frac{\mu_3^2}{2} > \frac{1}{2}(m_{FCNH})^2. \]

(B15)

(B16)

Several observations can be made from these relations: First is that for \( M_{W_R} \lesssim 6.5 \text{ TeV} \), one has \( \rho_{1,2,4} > 1 \). Thus, discovery of a \( W_R \) at the LHC would indicate a strongly coupled triplet sector. Second is that a small hierarchy among the \( \rho_i \) may exist. Requiring both \( H^0_{W_R} \) and \( A^0_{W_R} \) to be heavier than \( m_{FCNH} \) suggests \( \rho_3 \gtrsim 6\rho_1 \). Fig. 13 plots the values of \( \rho_3 \) for given \( M_{W_R} \) and \( m_{FCNH} \), and shows, for example, \( \rho_3 < 1 \) and \( m_{FCNH} \sim 15 \text{ (20) TeV} \) require \( M_{W_R} \gtrsim 10 \text{ (12) TeV} \). If \( H^0_{W_R} \) and \( A^0_{W_R} \) are largely responsible for neutral flavor transitions, then \( \rho_{1,3} \) can be reduced while keeping their differences fixed. We do not apply theoretical prejudices against strongly coupled systems and treat this as a consistent prediction. A more detailed discussion on the perturbativity of the scalar sector can be found in [52].

An ambiguity arises for the bidoublet self-couplings \( \lambda_{1,3,4} \) as the self-coupling of the SM-like Higgs is unconstrained. Using Eq. (B17), we take without impacting our study

\[ \lambda_1 \approx \frac{m_{H^0_{W_R}}^2}{2k_+}, \quad \lambda_2, \lambda_3 = 0. \]

(B17)
FIG. 13. Scalar triplet coupling $\rho_3$ contours for given $M_{W_R}$ and $m_{FCNH}$.

2. Configuration of LRSM FeynRules File

We implement our configuration of the scalar potential and choice for quark and lepton mixing into one FR restriction file that can be invoked when generating UFOs for the Manifest LRSM v1.1.6-MIX model file by \[73\]. See \[102\] for instructions. Internal parameters, e.g., $v_R$ and SM inputs of Eq. (28), can be modified via MG5 input parameter cards. The restriction file, lrsmLHCRestrictions.rst, is available from the source directory for the arXiv preprint version of this report \[127\]. It contains the following parameter identifications:

(* Turn off CKM mixing *)
\[s_{12} \rightarrow 0,\]
\[s_{23} \rightarrow 0,\]
\[s_{13} \rightarrow 0,\]

(* Turn off light neutrino mixing and set PMNS to diagonal *)
\[s_{L13} \rightarrow 0,\]
\[s_{L23} \rightarrow 0,\]
\[s_{L13} \rightarrow 0,\]

(* Turn off off-diagonal heavy/light neutrino mixing [V,X in Eq.(A.11) of 0901.3589]*)
\[V_{Ke} \rightarrow 0,\]
\[V_{Kmu} \rightarrow 0,\]
\[V_{Kta} \rightarrow 0,\]

(* Make mixing in LRSM manifest: all +1. Quasi-manifest: at least one -1 *)
\[W_{U11} \rightarrow 1,\]
\[W_{U22} \rightarrow 1,\]
\[W_{U33} \rightarrow 1,\]
\[W_{U11} \rightarrow 1,\]
\[W_{U22} \rightarrow 1,\]
\[W_{U33} \rightarrow 1,\]
\[W_{D11} \rightarrow 1,\]
\[W_{D22} \rightarrow 1,\]
\[W_{D33} \rightarrow 1,\]

(* LH vev*)
\[v_L \rightarrow 0,\]

(* Quark masses and Yukawas*)
\[M_U \rightarrow 0,\]
MD → 0,
MC → 0,
MS → 0,

(* Lepton masses and Yukawas *)
Me → 0,
Mμ → 0,
Mτ → 0,
MN1 → 0,
MN2 → 0,
MN3 → 0

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