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Abstract

We investigate left-right symmetric extensions of the standard model based on open strings ending on D-branes, with gauge bosons due to strings attached to stacks of D-branes and chiral matter due to strings stretching between intersecting D-branes. The left-handed and right-handed fermions transform as doublets under $Sp(1)_L$ and $Sp(1)_R$, and so their masses must be generated by the introduction of Higgs fields in a bi-fundamental $(\mathbf{2}, \mathbf{2})$ representation under the two $Sp(1)$ gauge groups. For such D-brane configurations the left-right symmetry must be broken by Higgs fields in the doublet representation of $Sp(1)_R$ and therefore Majorana mass terms are suppressed by some higher physics scale. The left-handed and right-handed neutrinos pair up to form Dirac fermions which control the decay widths of the right-handed W' boson to yield comparable branching fractions into dilepton and dijets channels. Using the most recent searches at LHC13 Run II with 2016 data we constrain the $(g_R, m_{W'})$ parameter space. Our analysis indicates that independent of the coupling strength g_R, gauge bosons with masses $m_{W'} \gtrsim 3.5$ TeV are not ruled out. As the LHC is just beginning to probe the TeV-scale, significant room for W' discovery remains.

I. INTRODUCTION

The $SU(3)_C \otimes SU(2)_L \otimes U(1)_Y$ standard model (SM) of particle physics has recently endured intensive scrutiny, with a dataset corresponding to an integrated luminosity of 12.9 fb^{-1} of 2016 pp collisions at $\sqrt{s} = 13$ TeV, and it has proven once again to be a remarkable structure that is consistent with all experimental results by tuning more or less 19 free parameters. However, the SM is inherently an incomplete theory, as it does not explain all known fundamental physical phenomena. The most obvious omission is that it does not provide a unification with gravity.

Superstring theory is the best candidate for a unified theory of all interactions. Dirichlet branes on which fundamental open string boundaries are attached are key objects to realize four-dimensional chiral gauge theories as low-energy effective models from superstring theory [1–4]. In general, gauge couplings are described by the string coupling, the Regge slope parameter $\alpha' = M_s^{-2}$, and partial volumes of compactified spaces. The conventional assumption is that M_s is of order (but below) the Planck scale, $M_{\rm Pl}$. Of particular interest is also the possibility of realistic superstring models with low mass string scale, $M_s \sim 10$ TeV, and large volume compactifications [5].

Intersecting D-brane models enclose a collection of building block guidelines, which can be used to manufacture the SM or something very close to it [6–14]. Within these models the basic unit of gauge invariance is a $U(1)$ field, and hence a stack of N identical D-branes sequentially gives rise to a $U(N)$ theory with the associated $U(N)$ gauge group. If there exist many types of D-brane, the gauge group grows into product form $\prod U(N_i)$, where N_i specifies the number of D-branes in each stack. (For $N = 2$, the gauge group can be $Sp(1) \cong SU(2)$ rather than $U(2)$.[1]) Gauge bosons (and associated gauginos in a supersymmetric model) arise from strings terminating on *one* stack of D-branes, whereas chiral matter fermions are realized as the Ramond-sector of open strings stretching between *two* stacks. For further details, see e.g. [15, 16].

The minimal embedding of the SM particle spectrum requires at least three brane stacks [17] leading to three distinct models of the type $U(3) \otimes U(2) \otimes U(1)$ that were classified in [17, 18]. Only one of them (model C of [18]) has baryon number as a gauge symmetry that guarantees proton stability (in perturbation theory), and can be used in the framework of low mass scale string compactifications. Besides, since the charge associated to the $U(1)$ of $U(2)$ does not participate in the hypercharge combination, $U(2)$ can be replaced by the symplectic $Sp(1)$ representation of Weinberg-Salam $SU(2)_L$, leading to a model with one extra $U(1)$ added to the hypercharge [19]. The SM embedding in four D-brane stacks leads to many more models that have been classified in [20, 21]. Several detailed and extensive phenomenological analyses have allowed us to blueprint new-physics signals of all these embeddings that could potentially be revealed at the LHC [22–35], in future e^+e^- and $\gamma\gamma$ colliders [36], as well as in astrophysical [37] and cosmological observations [31, 38, 39].

Curiously, the SM is chiral; it is of considerable interest to investigate ways to restore the left-right symmetry. In this paper we examine this possibility within the context of D-brane string compactifications. The layout of the paper is as follows. We begin in Sec. II with an outline of the basic setting of the minimal left-right symmetric intersecting D-brane model and a discussion on similarities and differences of the string inspired gauge structure and the canonical gauge sector generally used to restore the left-right symmetry [40–44]. Aspects of

[1] In the presence of orientifolds, one also obtains orthogonal and symplectic gauge groups.

coupling unification are briefly discussed in Sec. III. After that, in Sec. IV we examine the constraints placed on the parameter space of the D-brane set up by a broad range of LHC searches, taking into account the full set of relevant production and decay channels. Finally, we summarize our results and draw our conclusions in Sec. V.

II. LEFT-RIGHT SYMMETRY FROM INTERSECTING D-BRANES

The minimal left-right symmetric intersecting D-brane model is described by the gauge group $U(3)_C \otimes Sp(1)_L \otimes U(1)_L \otimes Sp(1)_R$ [10, 45]. The D-brane content is given in Table I; the mirror branes a^*, b^*, c^*, d^* are not shown. The right-handed quarks live at the intersection of the $Sp(1)_R$ stack of the D-branes and the color stack. The right-handed leptons stretch between $Sp(1)_R$ and the lepton brane $U(1)_L$. The left handed fermions bridge a mirror copy at the intersections of $U(3)_C \leftrightharpoons Sp(1)_L$ and $Sp(1)_L \leftrightharpoons U(1)_L$. The model requires two types of Higgs to generate the masses for the fermions and gauge bosons. The massless spectrum in the a, b, c, d basis is summarized in Table II. The baryon and lepton number, B and L, respectively; can be identified with the charges

$$B = \frac{1}{3} Q_a \qquad \text{and} \qquad L = Q_c \tag{1}$$

of the abelian gauge factors

$$U(1)_B = \frac{1}{3} U(1)_a \subset U(3)_C \qquad \text{and} \qquad U(1)_L = U(1)_d \,. \tag{2}$$

Note that both B and L are anomalous, but the combination $B - L$ is anomaly free.

Now a point worth noting at this juncture is that, in principle, one would like to introduce the left-right symmetry imposing a discrete $\mathbb{Z}_2$ invariance. This would be automatic if the gauge group could be $O(4)$, but then open strings would have no spinors and thus no doublets (under one of the $Sp(1)$) to break the symmetry. One should therefore impose a $\mathbb{Z}_2$ symmetry upon the exchange of the $Sp(1)_L$ and $Sp(1)_R$ branes, $\mathscr{P} : Sp(1)_L \leftrightharpoons Sp(1)_R$. This parity symmetry implies in particular equality of the two $Sp(1)$ gauge couplings and provides an interesting constraint for determining all couplings at the string scale. To keep our discussion general, in what follows we also consider models in which $\mathscr{P}$ is broken explicitly in this Higgs sector. These $\cancel{\mathscr{P}}$ models are peculiarly interesting as they could provide gauge coupling unification [45].

The first type of Higgs is in a bi-fundamental $(\mathbf{2}, \mathbf{2})$ representation under the two $Sp(1)$. The D-brane configuration may contain one or more bi-doublet scalars,

$$\Phi_i = \begin{pmatrix} \phi_{1i}^0 & \phi_{2i}^+ \\ \phi_{1i}^- & \phi_{2i}^0 \end{pmatrix}, \qquad \text{with} \qquad \tilde{\Phi}_i = \tau_2 \Phi_i^* \tau_2 \,, \tag{3}$$

to which correspond the covariant derivative

$$\mathcal{D}^\mu \Phi_i = \partial \Phi_i - i g_L \, \vec{W}_L^\mu \cdot \frac{\vec{\tau}}{2} \, \Phi_i + i g_R \, \Phi_i \, \frac{\vec{\tau}}{2} \cdot \vec{W}_R^\mu \,, \tag{4}$$

where g_L and g_R are gauge coupling constants, $\vec{W}_{L,R}^\mu$ are the $Sp(1)_{L,R}$ gauge fields, and $\vec{\tau}$ are the Pauli matrices. The bi-doublet scalars couple to fermion bilinear $\bar{Q}_L Q_R$ and $\bar{L}_L L_R$,

TABLE I: D-brane content of $U(3)_C \otimes Sp(1)_L \otimes U(1)_L \otimes Sp(1)_R$.

Label	Stack	Number of Branes	Gauge Group
a	Color	$N_a = 3$	$U(3)_C = SU(3)_C \times U(1)_a$
b	Left	$N_b = 1$	$Sp(1)_L \cong SU(2)_L$
c	Lepton	$N_c = 1$	$U(1)_L$
d	Right	$N_d = 1$	$Sp(1)_R \cong SU(2)_R$

TABLE II: Massless left-handed spectrum of $U(3)_C \otimes Sp(1)_L \otimes U(1)_L \otimes Sp(1)_R$.

Number	Fields	Sector	Representation	Q_a	Q_c	$B-L$
3	Q_L	$a \leftrightharpoons b$	$(\mathbf{3}, \mathbf{2}, \mathbf{1})$	1	0	1/3
3	$(Q_R)^c$	$a \leftrightharpoons d$	$(\bar{\mathbf{3}}, \mathbf{1}, \mathbf{2})$	-1	0	$-1/3$
3	L_L	$b \leftrightharpoons c$	$(\mathbf{1}, \mathbf{2}, \mathbf{1})$	0	1	-1
3	$(L_R)^c$	$c \leftrightharpoons d$	$(\mathbf{1}, \mathbf{1}, \mathbf{2})$	0	-1	1
N_Φ	Φ	$b \leftrightharpoons d$	$(\mathbf{1}, \mathbf{2}, \mathbf{2})$	0	0	0
N_{H_R}	H_R^u	$c \leftrightharpoons d$	$(\mathbf{1}, \mathbf{1}, \mathbf{2})$	0	1	-1
N_{H_R}	H_R^d	$c \leftrightharpoons d$	$(\mathbf{1}, \mathbf{1}, \mathbf{2})$	0	-1	1
N_{H_L}	H_L^u	$b \leftrightharpoons c$	$(\mathbf{1}, \mathbf{2}, \mathbf{1})$	0	-1	1
N_{H_L}	H_L^d	$b \leftrightharpoons c$	$(\mathbf{1}, \mathbf{2}, \mathbf{1})$	0	1	-1

and generate masses for quarks and leptons after spontaneous symmetry breaking by their vacuum expectation values (VEVs)

$$\langle \Phi_i \rangle = \frac{1}{\sqrt{2}} \, \mathrm{diag}(k_{1i}, k_{2i}) \,. \tag{5}$$

The left-right symmetry is broken by Higgs fields in the doublet representation of $Sp(1)_R$ with $B-L=1$. Note that we are forced to introduce these Higgs fields in vector-like pairs $H_R^{u,d}$ for anomaly cancellation. Here $\langle H_R^u \rangle = \binom{0}{v_1}$, $\langle H_R^d \rangle = \binom{v_2}{0}$, $v_R = \sqrt{v_1^2 + v_2^2} \gg k_{1i}, k_{2i}$ is $\mathcal{O}(\mathrm{TeV})$, and $\tan \varkappa \equiv v_1/v_2$. The associated left-handed Higgs doublets H_L^u and H_L^d must also be present to maintain the $\mathbb{Z}_2$ symmetry. Likewise, $\langle H_L^u \rangle = \binom{0}{v_3}$, $\langle H_L^d \rangle = \binom{v_4}{0}$, $v_L = \sqrt{{v_3}^2 + {v_4}^2}$, and $\tan \beta \equiv v_3/v_4$. In the analysis below we greatly simplify the discussion by assuming that H_L^u and H_L^d acquire vanishing, phenomenologically irrelevant, VEVs which are set to zero throughout; that is $v_L = 0$.

In the quark-mass basis the right-handed charged-current (CC) interaction of the $W_R^\pm$ boson for quarks is given by

$$\mathscr{L}_{\mathrm{CC}} = \frac{g_L}{\sqrt{2}} \overline{U}_L \gamma^\mu \mathbb{V}_L D_L W_{L\mu}^+ + \frac{g_R}{\sqrt{2}} \overline{U}_R \gamma^\mu \mathbb{V}_R D_R W_{R\mu}^+ + \mathrm{h.c.}\,, \tag{6}$$

where $\mathbb{V}_L$ is the SM CKM matrix in the canonical form, and $\mathbb{V}_R$ its right-handed equivalent, which has in principle different angles and five extra phases. The $\mathscr{P}$-symmetric case forces $|\mathbb{V}_L| = |\mathbb{V}_R|$, thus avoiding flavor-changing neutral currents. Since additional parameter freedom is currently superfluous, in all our calculations we assume the validity of the relation between the left- and right-handed CKM matrices imposed by the $\mathbb{Z}_2$ symmetry. More unconventional models have been discussed in e.g. [46–49]. It is noteworthy that the gauge fields $W^{\pm}_{L,R}$ are not quite mass eigenstates. The two charged gauge-bosons mix because both of them couple to the bi-doublet that is charged under the two $Sp(1)$ groups. The mass terms for the charged gauge-bosons are given by

$$\mathscr{L}_{m_W m_{W'}} = \begin{pmatrix} W^-_{L\mu} & W^-_{R\mu} \end{pmatrix} \begin{pmatrix} m_W^2 & \zeta m_W^2 \\ \zeta m_W^2 & m_{W'}^2 \end{pmatrix} \begin{pmatrix} W_L^{+\mu} \\ W_R^{+\mu} \end{pmatrix}, \tag{7}$$

where

$$m_W^2 = \frac{g_L^2}{4} \sum_i (k_{1i}^2 + k_{2i}^2) \qquad \text{and} \qquad m_{W'}^2 = \frac{g_R^2}{4} \left[2 v_R^2 + \sum_i (k_{1i}^2 + k_{2i}^2) \right], \tag{8}$$

and where the mixing is parameterized by an $\mathcal{O}(1)$ coefficient $\kappa = g_R/g_L$, with

$$\zeta = \kappa \, \frac{2 \sum_i k_{1i} k_{2i}}{\sum_i (k_{1i}^2 + k_{2i}^2)} \,. \tag{9}$$

The mass eigenstates $W_{1,2}$ (where $W_1 = W$ is identified as the well-known lighter state and $W_2 = W'$), are related to the gauge eigenstates W_L, W_R by a rotation of mixing angle ϕ, given by

$$\begin{pmatrix} W_L^{+\mu} \\ W_R^{+\mu} \end{pmatrix} = \begin{pmatrix} \cos\phi & -\sin\phi \\ \sin\phi & \cos\phi \end{pmatrix} \begin{pmatrix} W_1^{+\mu} \\ W_2^{+\mu} \end{pmatrix} \qquad \text{with} \qquad \tan 2\phi = \frac{-2\zeta m_W^2}{m_{W'}^2 - m_W^2} \ll 1 \,. \tag{10}$$

It is evident that in D-brane string compactifications the left-right symmetry cannot be broken by a Higgs field in the triplet representation of $Sp(1)_R$ with $B - L = \pm 2$, because the open string of such a massless mode would require four instead of two ends. As a consequence, there is no equivalent of the seesaw mechanism to generate the Weinberg term [50] which gives rise to Majorana neutrinos. To generate Majorana masses (by Q_c charge conservation) one needs dim-5 operators (such as $(L_R H_R^u)^2$) which are expected to be suppressed by the string scale. At the renormalizable level one can only write a coupling of leptons with the Higgs bi-doublet, which gives a Dirac mass to the neutrinos. As we show in the next section, the resulting neutrino mass constrains the W' decay channels, narrowing the parameter space for LHC searches.

The *physical* neutral gauge bosons Z_μ, Z'_μ and the photon A_μ are related to the weak $Sp(1)_{L,R}$ and $U(1)_{B-L}$ states $W^3_{R\mu}$, $W^3_{R\mu}$ and B_μ by an analogous mixing matrix. Using $v_R^2 \gg k_{1i}, k_{2i}$ the mass ratio of W' and Z' is found to be

$$\frac{m_{Z'}^2}{m_{W'}^2} \simeq \frac{\kappa^2 (1 - \sin^2\theta_W)}{\kappa^2 (1 - \sin^2\theta_W) - \sin^2\theta_W} > 1 \,, \tag{11}$$

with θ_W the Weinberg angle and $m_Z \simeq m_W / \cos\theta_W$ [51, 52].

In the foreground the chiral multiplets harbor a $[U(1)_a Sp(1)_L^2]$ mixed anomaly through triangle diagrams with fermions running in the loop. It is straightforward to see that the only anomaly free combination is $B-L$. The anomaly of the orthogonal combination is cancelled by the generalized Green-Schwarz mechanism, wherein closed string couplings yield classical gauge-variant terms whose gauge variation cancels the anomalous triangle diagrams [53–57]. The extra abelian gauge field becomes massive by the Green-Schwarz anomaly cancellation, behaving at low energies as a Z'' with a Stückelberg mass in general lower than the string scale by an order of magnitude, corresponding to a loop factor. Higgs VEVs will also generate additional mass terms for Z'', introducing also some small mixing with other gauge bosons, of order $(\text{TeV}/M_s)^2$. Note that for models with low mass string scale, the discovery of Z'' is within the LHC reach.

It is worth commenting on an aspect of this study which may seem discrepant at first blush. In principle, Euclidean brane instantons might contribute to Majorana masses [58, 59]. However, this would only be possible if the left-right breaking scale is of the order of the string scale, and consequently the W' and Z' gauge bosons would be out of the LHC reach. Hence, for a TeV-scale left-right symmetry breaking, we can generically argue as before that the strong suppression of Majorana mass terms in D-brane models constrains the W' decay channels, narrowing the parameter space for LHC searches.

III. GAUGE COUPLING UNIFICATION

For conventional models with 10 TeV $\ll M_s \lesssim M_{\rm Pl}$, there are still theoretical differences among the various frameworks which unfortunately do not easy lend themselves to experimental observation. In particular, it would be interesting to see what kind of unification scales (if any) are predicted by D-brane models in terms of running in the gauge sector.

To this end we can define $\alpha_i = g_i^2/(4\pi)$ as usual, and evolve α_i^{-1} which becomes simply $d\alpha_i^{-1}/dt = -b_i/2\pi$. The corresponding β-coefficients for the SM are found to be

$$(b_s^{\rm SM},\ b_L^{\rm SM},\ b_Y^{\rm SM}) = (-7,\ -19/6,\ 41/10)\,, \tag{12}$$

and those of the left-right symmetric extension are given by

$$(b_s,\ b_L,\ b_R,\ b_{B-L}) = (-7,\ -3,\ -3,\ 4) + (\delta b_s,\ \delta b_L,\ \delta b_R,\ \delta b_{B-L}), \tag{13}$$

where δb_i stands for the contributions from additional fields, not accounted for in the SM. For illustration, the coefficients for the groups $Sp(1)_L \otimes Sp(1)_R$ include the contribution from one bi-doublet field $\Phi_{\mathbf{1},\mathbf{2},\mathbf{2},0}$. We have included this field in the b_i directly, since the SM Higgs is $\Phi_{\mathbf{1},\mathbf{2},1/2} \in \Phi_{\mathbf{1},\mathbf{2},\mathbf{2},0}$ in our construction. At the left-right breaking scale $m_{W'}$ the hypercharge coupling splits into the $SU(2)_R$ and the $U(1)_{B-L}$ coupling according to

$$\frac{1}{\alpha_Y(m_{W'})} = \frac{3}{5}\,\frac{1}{\alpha_R(m_{W'})} + \frac{2}{5}\,\frac{1}{\alpha_{B-L}(m_{W'})}\,. \tag{14}$$

We can use (14) to reduce the system of equations by eliminating one of the four running gauge couplings, because the orthogonal combination $-2\alpha_R^{-1}(m_{W'})/5 + 3\alpha_{B-L}^{-1}(m_{W'})/5$ is a free parameter. Note that by gauge couplings we mean the independent parameters at the string scale, since $U(3)$ unifies the abelian g'_s with the non-abelian g_s with the appropriate normalisation: $g'_s(M_s) = g_s(M_s)/\sqrt{6}$ [17]. The abelian couplings of the D-brane model are furthered constrained by the orthogonality condition [30]. Finding a model which unifies correctly, then simply amounts to calculating a set of consistency conditions on the δb_i.

It is clear that by demanding $\mathbb{Z}_2$ symmetry we have $\delta b_R = \delta b_L$ and so, if we impose gauge coupling unification, (14) gives a left-right breaking scale out of the LHC reach. On the other hand, if $\delta b_R \neq \delta b_L$ then $m_{W'}$ could be $\mathcal{O}$(TeV). In addition, the restoration of the left-right symmetry can accelerate the running of the gauge couplings, yielding unification at a relatively low energy of about 10^{15} GeV, see e.g. [60]. Note that for the canonical left-right symmetric models, Super-K lower limits on half-life for proton decay [61–63] shift the scale of gauge coupling unification to higher energies [64]. However, this constraint does not apply to (left-right symmetric) intersecting D-brane models, which have baryon number as gauge symmetry that guarantees proton stability (in perturbation theory).

It is important to stress that for intersecting D-brane constructions, one can just require $g_L(M_s) = g_R(M_s)$ because of the $\mathbb{Z}_2$ symmetry (which makes the group $SO(4)$ effectively). The other gauge couplings are independent since they correspond to different brane stacks. This allows for an alternative to gauge coupling unification, with $m_{W'}$ of order TeV.

The role played by the renormalization-group flow in supersymmetric extensions of left-right symmetric D-brane models has been studied in [45, 65]. For $\mathscr{P}$ models, it is possible to get gauge coupling unification together with $m_{W'} \sim$ TeV. To accommodate the non-observation of supersymmetry signals with a TeV-scale left-right breaking we assume the following hierarchy of mass scales $m_{W'} < M_{\rm SUSY} < M_s$, where $M_{\rm SUSY}$ is the scale of supersymmetry breaking. For $N_{H_L} = 0$, this sets a lower bound on the number of vector-like pairs N_{H_R} for given a bi-doublet configuration N_Φ. For example, $N_\Phi = 1$ requires $N_{H_R} \geq 3$; otherwise $M_{\rm SUSY} < m_{W'}$ that is in conflict with our assumption. Following the extended survival hypothesis [66], in the energy regime $m_{W'} < E < M_{\rm SUSY}$ we take the minimal particle content of the non-supersymmetric left-right symmetric SM, that is one scalar Higgs bi-doublet Φ and one scalar Higgs doublet H_R^u, yielding

$$(\delta b_s,\ \delta b_L,\ \delta b_R,\ \delta b_{B-L}) = (0,\ 0,\ 1/6,\ 1/4)\ . \tag{15}$$

The β-function coefficients in the supersymmetric region are found to be

$$(\delta b_s,\ \delta b_L,\ \delta b_R,\ \delta b_{B-L}) = (4,\ 3 + N_\Phi,\ 3 + N_\Phi + N_{H_R},\ 2 + 3N_{H_R}/2)\ . \tag{16}$$

For $m_{W'} \sim 2$ TeV, the minimal possible scale of supersymmetry breaking comes out fairly universal $M_{\rm SUSY} \sim 19$ TeV; conjointly, the $Sp(1)_R$ gauge coupling does only vary slightly in the region $0.48 < g_R(m'_W) < 0.60$ [45].

Note that non-minimal D-brane constructs with 10 TeV $\ll M_s \lesssim M_{\rm Pl}$ could have more than one linear combination of anomalous $U(1)$ that are non-anomalous. Under certain topological conditions the associated gauge bosons can remain massless and obtain a low mass scale via the ordinary Higgs mechanism [67]. Interestingly, the extended scalar sector of these setups can be used to stabilize the vacuum up to the string scale [68]. Some phenomenological aspects of these kind of $U(1)$ gauge bosons and the prospects to search for them at the LHC were analyzed elsewhere [31].

As a matter of fact, the LHC8 phenomenology and discovery reach of massive Z' and Z'' gauge bosons have been discussed in detail in our previous publications [29–31]. The constraints on the parameter space are largely dominated by dijet searches. Since the dijet limits from LHC8 with an integrated luminosity of about 20 fb^{-1} [69, 70] are comparale to those of LHC13 with above about 12 fb^{-1} [71–73], in the next section we focus attention on W' searches.

IV. LHC PHENOMENOLOGY: CONSTRAINTS FROM W' SEARCHES

The ATLAS and CMS experiments are actively looking for W' and Z' gauge bosons. In particular, searches for W' resonances have been carried out at LHC8 and LHC13 considering a sequential SM W' [74] and the usual decay modes: the leptonic channels ($W' \to \tau\nu$ [75, 76], $W' \to \mu\nu$ [77], $W' \to e\nu$ [78], $W' \to l\nu$ [79–82]), the $W' \to WZ$ channel [83–87], the dijet [69–73] and the tb modes [88–93]. The sequential SM contains extra heavy neutral bosons Z' and W', with the same couplings to fermions and bosons as the Z and W. The limits derived by LHC experiments then apply directly to $\mathscr{P}$ models with $\mathbb{Z}_2$ symmetry. The lower bound on the mass of W' is 3.5 TeV at 95% C.L [72]. In this section we will extract from the results of the LHC analyses 95% C.L. exclusion regions on the $(m_{W'}, g_R)$ plane for $\not{\mathscr{P}}$ models, in which $g_R < g_L$.

The decay rate of W'^+ to fermion pairs is found to be

$$\Gamma(W'^+ \to u\bar{d}) = \Gamma(W'^- \to e^-\nu) = 3\kappa^2 A \left(1 + \frac{\alpha_s(m_{W'})}{\pi}\right) , \tag{17}$$

$$\Gamma(W'^+ \to t\bar{b}) = 3\kappa^2 A \left(1 + \frac{\alpha_s(m_{W'})}{\pi}\right)\left(1 - \frac{m_t^2}{m_{W'}^2}\right)^2 \left(1 + \frac{1}{2}\frac{m_t^2}{m_{W'}^2}\right) , \tag{18}$$

where $A = G_F m_W^2 m_{W'}/(6\pi\sqrt{2})$ is an overall constant [94, 95]. A first prediction of the D-brane model is then that W'^+ *cannot be leptophobic*, with comparable widths to dilepton and dijet channels.

The partial decay width to diboson is found to be

$$\begin{aligned}\Gamma(W'^+ \to W^+ Z) \;=\; & \frac{A}{4}\left(\frac{g_L^2}{g_Y^2 + g_L^2}\right)^2 a^2 \left(1 - 2\frac{m_W^2 + m_Z^2}{m_{W'}^2} + \frac{(m_W^2 - m_Z^2)^2}{m_{W'}^4}\right)^{3/2} \\ & \times \left(1 + 10\frac{m_W^2 + m_Z^2}{m_{W'}^2} + \frac{m_W^4 + 10 m_W^2 m_Z^2 + m_Z^4}{m_{W'}^4}\right) ,\end{aligned} \tag{19}$$

with $a = (\sin\phi/\kappa)(m_{W'}^2/m_W^2)$ [51]. The second prediction of the D-brane model is that (unless κ is small) *the decay rate into diboson is smaller than that into leptons:*

$$\frac{\Gamma(W'^+ \to W^+ Z)}{\Gamma(W'^+ \to ll)} \sim \frac{0.0121299}{\kappa^2} . \tag{20}$$

Note, however, that κ is bounded from below $\kappa > 0.55$ [51]. The reason is the following: the unbroken $U(1)$ is associated with $T_L^3 + T_R^3 + B - L$ and therefore (14). As a result g_R cannot be too small.

Following [96, 97], we compute the W' production cross section for pp collisions at $\sqrt{s} =$ 8 TeV multiplying the leading-order cross sections computed with MadGraph [98] (using model files generated with FeynRules [99] and CTEQ6L parton distributions [100], with factorization and renormalization scales set at $m_{W'}$) by a scale-dependent K-factor which takes into account next-to-leading order (NLO) QCD effects. We adopt $1.32 \lesssim K \lesssim 1.37$ as derived in [101], which is larger than both the K-factor obtained with the parton level Monte Carlo program for FeMtobarrn processes (MCFM) [102] and the factor $K \approx 1.15$ computed in [103]. With this in mind, we parametrize the W' production cross section at LHC8 by

$$\sigma(pp \to W')_{\sqrt{s}=8~\mathrm{TeV}} = 816.686\; g_R^2\; \exp\left[-3.53131\left(\frac{m_{W'}}{\mathrm{TeV}}\right)\right]\;\mathrm{pb}\,. \tag{21}$$

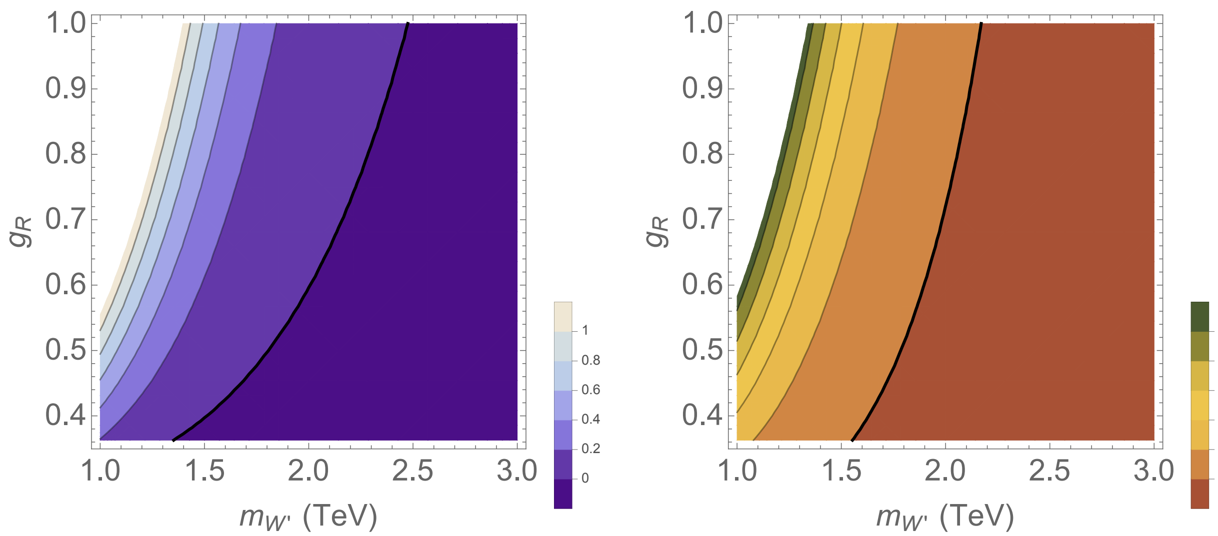

FIG. 1: Contours of constant $\Delta\sigma_X$ for $X =$ dijet (left) and tb (right) modes for a generous range of g_R, and LHC8 collisions.

We note that, for $m_{W'} = 2$ TeV, the value of the parametrization is consistent with the results of [104]. However, the results of our parametrization are in general a factor of ≈ 8 larger than the parametrization given in [52]. We note that when we multiply (21) by (18) we recover the inclusive cross section for tb modes derived in [103].

For pp collisions at $\sqrt{s} = 13$ TeV, we fit the results derived in [105] for the W' production cross section at NLO with threshold resummation at next-to-next-to-leading logarithm (NNLL) matched to threshold-improved parton distributions functions. We parametrize the W' production cross section at LHC13 by

$$\sigma(pp \to W')_{\sqrt{s}=13 \text{ TeV}} \simeq 3136.28 \; g_R^2 \; \text{AntiLog}_{10}[f(m_{W'})] \text{ pb}\,, \tag{22}$$

where

$$f(m_{W'}) = -1.61679 \left(\frac{m_{W'}}{\text{TeV}}\right) + 0.10953 \left(\frac{m_{W'}}{\text{TeV}}\right)^2 - 0.00385164 \left(\frac{m_{W'}}{\text{TeV}}\right)^3 . \tag{23}$$

We now turn to bound the parameter space of the D-brane model. To do so, we define

$$\Delta\sigma_X = \sigma(pp \to W') \times \mathcal{B}(W' \to X) - [\sigma(pp \to W') \times \mathcal{B}(W' \to X)]_{95\%\text{CL}}\ , \tag{24}$$

where $[\sigma(pp \to W') \times \mathcal{B}(W' \to X)]_{95\%\text{CL}}$ is the 95% C.L. upper limit on the inclusive cross section $\sigma(pp \to W') \times \mathcal{B}(W' \to X)$, with $X =$ dijet and tb modes. Our results are encapsulated in Figs. 1 and 2 where we show contours of constant $\Delta\sigma_X$ for pp collisions at $\sqrt{s} = 8$ TeV and $\sqrt{s} = 13$ TeV. The dark line corresponds to $\Delta\sigma_X = 0$, with the parameter space to the left being excluded at the 95% C.L. In the plot we also show contours to the left of the exclusion line in order to leave room to move the bound as necessary to account for detector effects and possible future change in theoretical uncertainties. Note that the exclusion line is conservative since our calculation is at the parton level. Therefore, taking into account the aforementioned effects would move the line to lower masses. As can be seen in Figs. 1 and 2, for the interesting range $0.48 < g_R(m'_W) < 0.60$, gauge bosons with $m_{W'} \gtrsim 2.5$ TeV are not excluded. Note that at present the most restrictive bounds on W' masses are from dijet searches, but an actual discovery will perhaps be possible using the tb modes, which have comparable sensitivity.

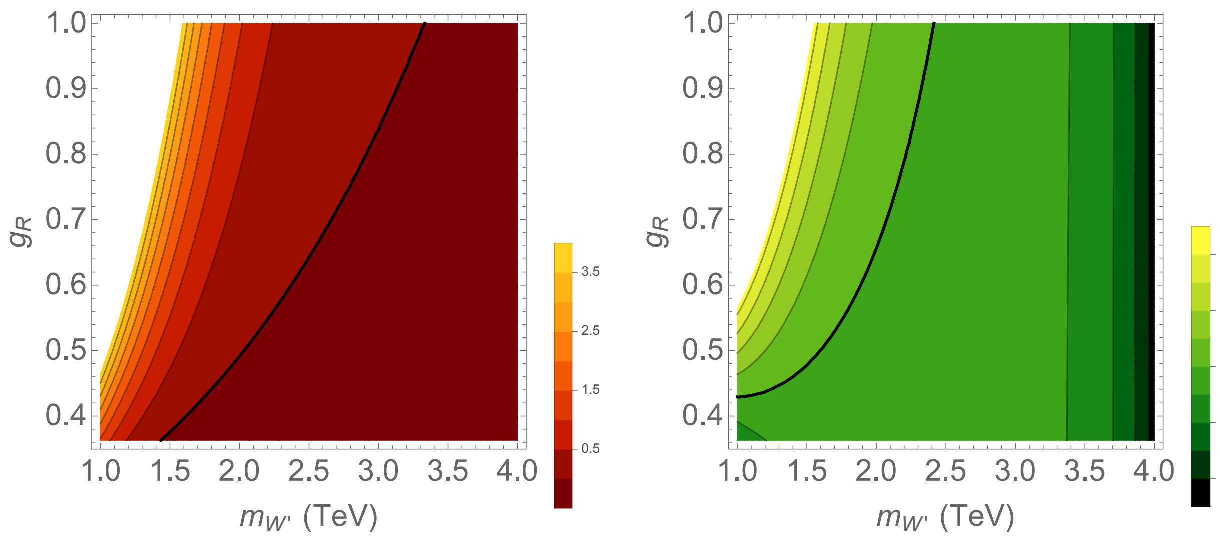

FIG. 2: Contours of constant $\Delta\sigma_X$ for X = dijet (left) and tb (right) modes for a generous range of g_R, and LHC13 collisions.

V. CONCLUSIONS

In this paper we studied extensions of the SM in which restoration of the left-right symmetry is pulled down to the TeV-scale ballpark. We considered string theory setups in which the gauge theories live on D$(3+p)$-branes which entirely fill the uncompactified part of space-time and wrap certain p-cycles inside the compact six-dimensional manifold. The chiral matter fields appear at the intersection of two D$(3+p)$-branes. To develop our program in the simplest way, we worked within the construct of a minimal model characterized by $U(3)_C \otimes Sp(1)_L \otimes U(1)_L \otimes Sp(1)_R$.

The approach we have taken in this work can be regarded as an effective theory with both interesting LHC phenomenology and unique theoretical characteristics such as conservation of B to prevent proton decay to all orders in perturbation theory and violation of L without Majorana masses. The absence of Majorana masses constrains the W' decay rates, yielding comparable widths to dilepton and dijet channels. This naturally narrows the parameter space for LHC searches while providing at the same time a definite prediction of D-brane models. Additional abelian gauge symmetries, inherent to the structure of D-brane gauge theories, can provide interesting corroboration for string physics near the TeV-scale.

We have derived bounds on the $(g_R, m_{W'})$ plane using the most recent searches at LHC13 Run II with 2016 data. Our analysis indicates that independent of the gauge coupling g_R, right-handed W' bosons with masses above about 3.5 TeV are not ruled out. As the LHC is just beginning to probe the TeV-scale, significant room for discovery remains.

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