Gauged baryon and lepton number in MSSM_{4} brane worlds

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Abstract

A recent D-brane model designed to accommodate a phenomenologically acceptable fourth generation of chiral fermions was noted to produce an unexpected additional unbroken nonanomalous U(1) gauge group at the string scale. We show that the corresponding charges acting on MSSM fields count baryon and lepton numbers. If broken spontaneously at lower scales, these U(1)$_B$ and U(1)$_L$ symmetries provide potential avenues for preserving baryogenesis while nonetheless explaining the suppression of proton decay (without the need for R parity), as well as the smallness of right-handed neutrino Majorana masses compared to the string scale.

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I. INTRODUCTION

The long lifetime of the proton, in excess of 10^{34} y, provides one of the strongest hints of fundamental physics beyond the electroweak scale. Baryon (B) and lepton (L) number are conserved as global symmetries of the Standard Model (SM) and the Minimal Supersymmetric Standard Model (MSSM). However, global symmetries in a theory can easily appear by accident, and unless they are protected in some way, are broken at some level by nonperturbative effects. Indeed, one of the distinctive features of grand unified theories (GUTs) is their introduction of new particles that can mediate proton decay. A typical non-supersymmetric GUT with arbitrary couplings of natural size allows proton decay at an unacceptably high rate, which [in addition to the taming of the gauge hierarchy problem and the numerical improvement of $\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ gauge unification] provides impetus for the study of supersymmetric (SUSY) theories. In particular, the scale at which unification occurs in SUSY models is pushed to a slightly higher energy, leading to an additional suppression of the proton decay rate.

However, unmodified SUSY theories generate their own problems with B and L nonconservation. While B- and L-violating terms in non-SUSY theories first appear in operators of dimension $d = 6$, the presence of squark and slepton fields allows such terms to arise already at $d = 4$. Denoting the left-handed chiral matter and Higgs superfields by the quantum numbers of B, L, and $\mathrm{SU}(3)_C \times \mathrm{SU}(2)_L \times \mathrm{U}(1)_Y$ in the usual way,

$$Q(+\tfrac{1}{3}, 0, \mathbf{3}, \mathbf{2}, +\tfrac{1}{6}), \;\; L(0, +1, \mathbf{1}, \mathbf{2}, -\tfrac{1}{2}), \;\; U^c(-\tfrac{1}{3}, 0, \overline{\mathbf{3}}, \mathbf{1}, -\tfrac{2}{3}), \;\; D^c(-\tfrac{1}{3}, -1, \overline{\mathbf{3}}, \mathbf{1}, +\tfrac{1}{3}),$$

$$N^c(-1, 0, \mathbf{1}, \mathbf{1}, 0), \;\; E^c(-1, 0, \mathbf{1}, \mathbf{1}, +1), \;\; H_u(0, 0, \mathbf{1}, \mathbf{2}, +\tfrac{1}{2}), \;\; H_d(0, 0, \mathbf{1}, \mathbf{2}, -\tfrac{1}{2}), \quad (1.1)$$

one finds already at $d = 4$ several B- and L-violating operators, such as $N^c N^c N^c$ and QD^cL (where the necessary color and weak isospin contractions are implied). Conventional R parity [1–5], which may be defined operationally on MSSM fields as $(-1)^{3(B-L)+2J}$, eliminates such terms while protecting the conventional Yukawa terms such as LH_uN^c. Nevertheless, even with this restriction Refs. [3, 6] showed that $d = 5$ operators such as $QQQL$ still lead to an unacceptably short proton lifetime unless additional suppressions are introduced.

A classic study [7] of possible B- and L-violating operators at $d = 4$ and 5 and how various discrete symmetries may be used to suppress phenomenologically unacceptable ("dangerous") decay rates also summarizes which of these operators are especially dangerous by

providing estimates of their suppressed coefficients. While the number of such operators increases greatly after allowing the right-handed neutrino field N^c [8] (which was not considered in [7]; we count 36 B- and/or L-violating operators up to and including $d = 5$), one can check that none of them are singlets under a symmetry that assigns a charge $Q_X = -1$ to Q and L, $+1$ to U^c, D^c, N^c, and E^c, and 0 to $H_{u,d}$ [for example, such a symmetry can occur as a subgroup of a hypothetical U(1)$_X$, as seen below]. On the other hand, the μ term $H_u H_d$ and the standard Yukawa terms are all singlets under this symmetry, indicating a potential route for suppressing proton decay without recourse to R parity.

Going beyond global symmetries, one may consider gauging the symmetries U(1)$_B$ or U(1)$_L$ that count baryon and lepton numbers, respectively, either for its own phenomenological interest in collider physics [9], or as a method of providing interesting dark matter candidates in both non-SUSY [10] and SUSY [11] models. So-called leptophilic matter [having a U(1) charge that couples preferentially to leptons rather than quarks, but not necessarily just U(1)$_L$] can serve to explain [12] the observed PAMELA [13]/ATIC [14]/Fermi LAT [15] cosmic ray positron excess, while a relatively light leptophobic [coupling primarily to quarks rather than leptons, but not necessarily just U(1)$_B$] Z' can serve to explain [16] Tevatron anomalies in the measured $t\bar{t}$ forward-backward asymmetry [17] and the associated production of Ws with jets [18]. Introducing additional Abelian symmetries beyond those of the SM or MSSM requires the consideration of a variety of new potential anomalies and the addition of new chiral matter multiplets to enforce their cancellation; addressing such concerns accounts for much of the analysis in the models mentioned above.

The introduction of almost any extra gauged U(1) that contains a portion coupling universally to quarks [$i.e.$, U(1)$_B$] leads to the prohibition of proton decay at energy scales such that the symmetry remains unbroken. Since the proton is the lightest fermionic hadron, its decay products must contain an odd number of leptons, and some potential combinations including U(1)$_B$, such as U(1)$_{B-L}$ (if present), would have to break at a somewhat higher scale than others in order to support the experimental lifetime limit. Indeed, such possibilities have been considered for decades [19]. However, demanding baryon number to be an exactly conserved quantity at all scales would also eliminate the possibility of baryogenesis and a natural resolution of the problem of matter-antimatter asymmetry in the universe.

Analogous to the possibility that unconstrained operators can generate excessively large proton decay but their complete elimination forbids mechanisms for generating baryogenesis,

the μ-term coupling also introduces a well-known difficulty [1] in the MSSM: The natural Planck- or string-scale value of μ would force comparable values for the Higgs masses, and therefore must be suppressed. However, if μ is taken too small, light charginos appear in mass ranges that have already been excluded by collider searches. In the original solution [1], SUSY is explicitly but softly broken by the μ term in the TeV range; however, one may also consider the possibility of μ arising from spontaneous symmetry breaking. An efficient model of beyond-MSSM physics addresses both issues, as in Refs. [8, 20]. Indeed, both of these papers also note (as suggested above) a diminished significance for R parity. Moreover, another prime motivation for R parity—the existence of a lightest SUSY particle as a dark matter candidate—is ameliorated if one has recourse to other selection rules based upon additional U(1) symmetries to produce other long-lived exotic particles.

Under what circumstances can one develop theories with nonanomalous U(1)s that help to stabilize the proton? Here too, one can find interesting observations in the literature. As noted in Refs. [21, 22], the desired U(1)s do not appear to fit into conventional SO(10) or E_6-type grand unification schemes. Amusingly, the former paper (by Pati) advocates a stringy origin for the extra U(1), while the latter paper (inspired by results of explicit models [23, 24]) advocates a Pati-Salam (left-right symmetric) scheme [25] as the origin of the extra U(1). The model we present here, as it turns out, supports both stringy and Pati-Salam ideas. The possibility of suppressing proton decay using a left-right symmetric heterotic string-derived U(1) is considered in Refs. [26, 27], and a fourth generation in heterotic string theory is considered in Refs. [28, 29].

Furthermore, cancelling the anomalies of the U(1) associated with the suppression of proton decay appears to be more natural in models that introduce an additional set of chiral fermions; Ref. [20] includes a full family of fermions with opposite chiralities to those of the SM, but also notes that similar cancellations are possible in a four-generation MSSM. The potential significance of four generations in this context is also explored in Refs. [30, 31]. This is precisely the scenario we propose here. While long considered phenomenologically unlikely, the existence of a complete fourth generation of light chiral fermions in the range of 100–600 GeV to be detected at the Large Hadron Collider (LHC) remains a very real possibility, an idea recently revived in Ref. [32]. As one might expect, the masses and mixings of such new fermions must satisfy a variety of constraints arising from direct searches, electroweak precision constraints, and perturbative unitarity, but substantial regions in parameter space

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remain for their discovery, all accessible to the LHC.

In light of these phenomenological possibilities, in recent papers [33, 34] we have developed two 4-generation models in Type IIA string theory using D6-branes intersecting on a $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold. Our interest in both cases was to explain the hierarchy of masses and mixings among the known fermions. In the first model [33], we obtained full rank-4 Yukawa matrices at the level of trilinear couplings, but fitting the large hierarchies in masses and mixings required the adjustable vacuum expectation values (VEVs) and open-string moduli to lie in very particular regions of parameter space. The second model [34], which is of direct interest for this paper, generates Yukawa matrices only of rank 2 at trilinear order: Only the heaviest two generations obtain masses at this order, while the lighter two generations obtain masses via higher-order effects, which arguably provides a better physical explanation for facts such as, *e.g.*, $m_t/m_e \approx 340{,}000$. Moreover, natural solutions simultaneously accommodating the known third-generation fermion masses and the constraints on the fourth-generation fermion masses are relatively easy to obtain in the model of [34].

For the reader unfamiliar with the details of building consistent string theory-based models containing the MSSM, the essential points are to obtain a compactification of the extra 6 spacetime dimensions that supports just the low-energy particle content of the theory and to develop mechanisms for eliminating any unwanted states arising from residual degrees of freedom. The gauge-group charges and multiplicities are determined by the geometry (angles and numbers of intersections in Type IIA) through which strings wrap the compactified dimensions. The requirement of SUSY, the elimination of anomalies, and the imposition of the Green-Schwarz mechanism [35] to allow for massless U(1) gauge symmetries (all essential to the topic of this paper) may all be expressed using simple well-known constraints in terms of the wrapping numbers, brane multiplicities, and form of the compactified space, particularly for the simple geometry of the 6-torus orientifold $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$.

As mentioned above, our construction uses a Pati-Salam (PS) SU(4)$_C$×SU(2)$_L$×SU(2)$_R$ scheme to obtain fermions of the correct charges: Branes a, b, and c carry color/lepton, right-chiral, and left-chiral charges, respectively. Unification of the low-energy MSSM gauge group SU(3)$_C$×SU(2)$_L$×U(1)$_Y$ occurs in an efficient and natural way, which in turn is made possible because U(1)$_Y$ remains massless and anomaly free by surviving the GS mechanism. However, the two models [33, 34] differ in one very distinctive manner: In the former, the only such nonanomalous U(1)s at the string scale (as is typical to PS constructions) are

$U(1)_{B-L}$ and $U(1)_{I_{3R}}$, from which $U(1)_Y = \frac{1}{2}U(1)_{B-L}+U(1)_{I_{3R}}$. However, the latter model generates an additional nonanomalous $U(1)_X$ coupling to the MSSM fields with precisely the charges described above. The crucial observation is that the combinations

$$U(1)_B \equiv -\frac{1}{4}U(1)_X + \frac{1}{4}U(1)_{B-L}, \; U(1)_L \equiv -\frac{1}{4}U(1)_X - \frac{3}{4}U(1)_{B-L} \qquad (1.2)$$

are nonanomalous U(1)s surviving at the string scale that, when acting on MSSM fields, precisely count baryon and lepton number, respectively. These remarkable features appear due to a confluence of a PS construction, the presence of four chiral fermion generations, and trilinear Yukawa matrices of less-than-full rank.

In this paper we consider the broad phenomenological consequences of the enhanced symmetries of this model. In Sec. II we exhibit beyond-MSSM operators that can contribute to proton decay, reprise the field content of the model of Ref. [34], and investigate which operators survive due to $U(1)_X$ conservation, while we consider the nature of the new symmetry in Sec. III. The phenomenology induced by the new symmetry is discussed in Sec IV: suppression of proton decay, the μ term, the right-handed neutrino Majorana mass term, and baryogenesis. Section V emphasizes the prospects for new physics in these models, and Sec. VI offers concluding thoughts.

II. PROTON STABILITY AND AN EXTRA GAUGED U(1)

Although supersymmetry provides an elegant solution to the gauge hierarchy problem, it introduces some additional complications. First among these is the rapid decay of the proton through the pair of $d = 4$ F-term operators (B- and L-violating, respectively):

$$U^c D^c D^c, \; QD^c L. \qquad (2.1)$$

This problem is usually solved in the MSSM by introducing R parity, under which the known fermions are even while their SUSY partners are odd (or the related "matter parity", under which $R = +1$ for Q, U^c, D^c, L, E^c, N^c and $R = -1$ for $H_{u,d}$). As a bonus, R parity leads to a stable lightest SUSY particle (LSP), which is a natural candidate for dark matter. Although this idea is attractive, it is well known that a gauged $U(1)_{B-L}$ also forbids the $d = 4$ operators, and furthermore, R parity [more specifically, matter parity $(-1)^{3(B-L)}$] can result from $U(1)_{B-L}$ broken spontaneously to its discrete $\mathbb{Z}_2$ subgroup [19, 36, 37].

Even though the problem of rapid proton decay via $d = 4$ operators can be eliminated through this mechanism, one still faces the problem of $d = 5$ operators that allow for proton decay with a lifetime too short to evade current experimental constraints unless the coefficients of these operators are chosen to be sufficiently small. First among these are single operators that allow (at least in principle) proton decay and preserve $B - L$:

$$[QQQL]_F \,, \ [U^c U^c D^c E^c]_F \,, \ [D^c D^c U^c N^c]_F \,. \tag{2.2}$$

The second set consists of relevant $d = 5$ operators that violate either B or L separately, which combine with the appropriate member of Eq. (2.1) to form a composite operator that conserves $B - L$ and allows proton decay:

$$[QQQH_d]_F \,, \ [QU^c E^c H_d]_F \,, \ [QU^c L^\dagger]_D \,, \ [U^c (D^c)^\dagger E^c]_D \,, \ [QQ(D^c)^\dagger]_D \,,$$
$$[QQ^\dagger N^c]_D \,, \ [U^c (U^c)^\dagger N^c]_D \,, \ [D^c (D^c)^\dagger N^c]_D \,, \ [QU^c N^c H_u]_F \,, \ [QD^c N^c H_d]_F \,. \tag{2.3}$$

Indeed, the $d = 5$ operators are the ones that effectively lead to the exclusion of GUTs based on minimal SU(5) [38], although these operators can be suppressed in other unified models, *e.g.*, flipped SU(5) [39, 40].

Let us consider the model constructed in Type IIA string theory with intersecting D6-branes discussed in [34]. At the string scale, it is a four-generation Pati-Salam (PS) model, with the gauge symmetry of the "observable" sector given by

$$\mathrm{U}(4)_C \times \mathrm{U}(2)_L \times \mathrm{U}(2)_R \,, \tag{2.4}$$

in addition to a "hidden" sector with the gauge group $\mathrm{USp}(8)^3$. Since $\mathrm{U}(N) = \mathrm{SU}(N) \times \mathrm{U}(1)$, each observable gauge group has an associated anomalous U(1), denoted here as $\mathrm{U}(1)_4$, $\mathrm{U}(1)_{2L}$, and $\mathrm{U}(1)_{2R}$, respectively (which correspond to the a, b, and c brane stacks, respectively). The anomalies of these Abelian symmetries are cancelled by a generalized Green-Schwarz mechanism, and as a result their gauge bosons obtain string-scale masses. However, these U(1)s remain as global symmetries to all orders in perturbation theory, and as a result perturbatively forbid operators that would otherwise be allowed.

As noted in [34], precisely one linear combination of $\mathrm{U}(1)_4$, $\mathrm{U}(1)_{2L}$, and $\mathrm{U}(1)_{2R}$ remains massless and anomaly free, which is given by

$$\mathrm{U}(1)_X \equiv \mathrm{U}(1)_4 + 2 \left[\mathrm{U}(1)_{2L} + \mathrm{U}(1)_{2R} \right] \,. \tag{2.5}$$

TABLE I: The chiral and vectorlike superfields, and their multiplicities and quantum numbers under the gauge symmetry $U(4)_C \times U(2)_L \times U(2)_R \times [USp(8)_1 \times USp(8)_2 \times USp(8)_3]$. The $U(1)_X$ charge Q_X is given by the combination $Q_X = Q_4 + 2(Q_{2L} + Q_{2R})$.

	Mult.	Quantum Number	Q_4	Q_{2L}	Q_{2R}	Q_X	Field
ab	4	$(4, \overline{2}, 1, 1, 1, 1)$	1	-1	0	-1	$F_L(Q_L, L_L)$
ac	4	$(\overline{4}, 1, 2, 1, 1, 1)$	-1	0	1	1	$F_R(Q_R, L_R)$
bc	4	$(1, 2, \overline{2}, 1, 1, 1)$	0	1	-1	0	H_u^i, H_d^i
$a1$	1	$(4, 1, 1, \overline{8}, 1, 1)$	1	0	0	1	X_{a1}
$a2$	1	$(\overline{4}, 1, 1, 1, 8, 1)$	-1	0	0	-1	X_{a2}
$b2$	1	$(1, 2, 1, 1, \overline{8}, 1)$	0	1	0	2	X_{b2}
$b3$	2	$(1, \overline{2}, 1, 1, 1, 8)$	0	-1	0	-2	X_{b3}^i
$c1$	1	$(1, 1, \overline{2}, 8, 1, 1)$	0	0	-1	-2	X_{c1}
$c3$	2	$(1, 1, 2, 1, 1, \overline{8})$	0	0	1	2	X_{c3}^i
b_S	2	$(1, 3, 1, 1, 1, 1)$	0	2	0	4	T_L^i
b_A	2	$(1, \overline{1}, 1, 1, 1, 1)$	0	-2	0	-4	S_L^i
c_S	2	$(1, 1, \overline{3}, 1, 1, 1)$	0	0	-2	-4	T_R^i
c_A	2	$(1, 1, 1, 1, 1, 1)$	0	0	2	4	S_R^i
ab'	2	$(4, 2, 1, 1, 1, 1)$	1	1	0	3	Ω_L^i
	2	$(\overline{4}, \overline{2}, 1, 1, 1, 1)$	-1	-1	0	-3	$\overline{\Omega}_L^i$
ac'	2	$(4, 1, 2, 1, 1, 1)$	1	0	1	3	Φ_i
	2	$(\overline{4}, 1, \overline{2}, 1, 1, 1)$	-1	0	-1	-3	$\overline{\Phi}_i$
bb'	4	$(1, \overline{1}, 1, 1, 1, 1)$	0	-2	0	-4	s_L^i
	4	$(1, 1, 1, 1, 1, 1)$	0	2	0	4	$\bar{s}_L^i$
cc'	4	$(1, 1, 1, 1, 1, 1)$	0	0	2	4	s_R^i
	4	$(1, 1, \overline{1}, 1, 1, 1)$	0	0	-2	-4	$\bar{s}_R^i$
bc'	1	$(1, 2, 2, 1, 1, 1)$	0	1	1	4	H'
	1	$(1, \overline{2}, \overline{2}, 1, 1, 1)$	0	-1	-1	-4	$\overline{H}'$

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Thus, the full gauge symmetry of the model at this stage is given by

$$\mathrm{SU}(4)_C \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R \times \mathrm{U}(1)_X \times \left[\mathrm{USp}(8)^3\right]. \tag{2.6}$$

We exhibit the matter spectrum and the quantum numbers of each field in Table I. As seen there, the superfields $F_L^i(Q_L, L_L)$ carry charge $\mathrm{U}(1)_X$ charge $Q_X = -1$, $F_R^i(Q_R, L_R)$ $(= U^c, D^c, N^c, E^c)$ carry charge $Q_X = +1$, and the Higgs superfields are uncharged under $\mathrm{U}(1)_X$, as promised in the Introduction. Thus, the Yukawa couplings

$$
\begin{aligned}
W_Y &= y_{ijk} F_L^i F_R^j H^k \tag{2.7} \\
&= y_{ijk}^u Q^i (U^c)^j H_u^k + y_{ijk}^d Q^i (D^c)^j H_d^k + y_{ijk}^\nu L^i (N^c)^j H_u^k + y_{ijk}^l L^i (E^c)^j H_d^k,
\end{aligned}
$$

where $i, j, k = \{1, 2, 3, 4\}$, are allowed by both the global $\mathrm{U}(1)$ symmetries as well as the gauged $\mathrm{U}(1)_X$ symmetry. As shown in [34], the resulting Yukawa matrices are rank 2, which allows for fermion mass textures that can easily satisfy constraints placed on fourth-generation fermion masses. Since the fields F_L and F_R have different $\mathrm{U}(1)_X$ charges, these fields obviously cannot be placed in the same spinorial **16** of SO(10). Thus, this PS model cannot have an SO(10) origin.

As can also easily be checked using Table I, none of the $d = 4$ and $d = 5$ operators given in Eqs. (2.1), (2.2), and (2.3) are singlets under $\mathrm{U}(1)_X$ (Nor, it should be added, are any of the remaining 21 operators of the set of 36 mentioned in the Introduction, all of which are purely L-violating). Indeed, since $\mathrm{U}(1)_X(Q) = +1$, $\mathrm{U}(1)_X(U^c, D^c) = -1$, $\mathrm{U}(1)_X(H_u, H_d) = 0$, one finds that it is not possible to create a B-violating $\mathrm{U}(1)_X$-singlet $d \leq 6$ operator without including exotic matter. Thus, B- and L-violating processes that proceed via these operators are effectively forbidden at scales where $\mathrm{U}(1)_X$ remains unbroken.

As mentioned in the previous section, the MSSM superfields are also charged under global symmetries that arise from anomalous $\mathrm{U}(1)$s whose gauge bosons obtain masses via the generalized Green-Schwarz anomaly cancellation mechanism. As can be seen from Table I, none of the dimension-4 and -5 operators that mediate proton decay are singlets under these global symmetries, and are thus perturbatively forbidden. However, these global symmetries can, in principle, be broken by nonperturbative effects, namely D-brane instantons [41–43], assuming that suitable instantons with the correct zero-mode structure and satisfying all constraints can actually arise.

If such suitable instantons are present in the model, then proton decay can be allowed via these operators, albeit at a rate exponentially suppressed by the instanton action. While it

might be possible for the instanton suppression to be large enough to allow proton decay at an acceptable rate, this scenario is by no means guaranteed. However, for the present model, these operators cannot be induced via D-brane instantons since the $U(1)_X$ gauge symmetry remains unbroken at the string scale. Therefore, proton decay via these operators is forbidden at scales where $U(1)_X$ remains unbroken, and highly suppressed at scales at which $U(1)_X$ is broken since the effects by which they may appear are perturbatively forbidden.

III.   THE $U(1)_L$- AND $U(1)_B$-EXTENDED MSSM'S

The PS gauge symmetry may be broken by a process of D6-brane splitting [44, 45], which corresponds to assigning VEVs to some of the adjoint scalars associated with each stack along its flat directions (See [34] for a more detailed discussion of the process for this model). For present purposes, we note that the splitting results in additional massless and anomaly-free $U(1)$s given by

$$
\begin{aligned}
U(1)_{B-L} &= \frac{1}{3}U(1)_{\text{baryon}} - U(1)_{\text{lepton}} , \\
U(1)_{I_{3R}} &= \frac{1}{2}U(1)_{\text{up}} - \frac{1}{2}U(1)_{\text{down}} ,
\end{aligned}
\tag{3.1}
$$

while $U(1)_X$ becomes

$$
\begin{aligned}
U(1)_X &= U(1)_{\text{baryon}} + U(1)_{\text{lepton}} + 2\left[U(1)_{2L} + U(1)_{\text{up}} + U(1)_{\text{down}}\right] \\
&= U(1)_{a1} + U(1)_{a2} + 2\left[U(1)_b + U(1)_{c1} + U(1)_{c2}\right] ,
\end{aligned}
\tag{3.2}
$$

where $U(1)_{\text{baryon}}=U(1)_{a1}$ and $U(1)_{\text{lepton}}=U(1)_{a2}$ are global $U(1)$s obtained from $U(1)_4 =U(1)_a$ after brane splitting, under which baryons and leptons are charged, respectively. Similarly, $U(1)_{\text{up}}=U(1)_{c1}$ is a global symmetry under which up-type quarks and neutrinos are charged, while down-type quarks and electrically-charged leptons are charged under the global symmetry $U(1)_{\text{down}}=U(1)_{c2}$, both obtained from $U(1)_{2R} =U(1)_b$ after brane splitting. Typically in this type of model, the charges of the quarks and leptons under these global (and indeed, anomalous) symmetries are identified as $\frac{1}{3}U(1)_{\text{baryon}}$ and $U(1)_{\text{lepton}}$, respectively. However, as we now show, in this model one finds it is possible to obtain *gauged* $U(1)$s under which the quarks and leptons have numerically identical charges, that we label as the true baryon number (B) and lepton number (L), respectively. In particular, when acting upon MSSM fields, the $U(1)_X$ gauge symmetry is effectively $-U(1)_{3B+L}$.

After brane splitting, the observable gauge symmetry of the model becomes $SU(3)_C \times SU(2)_L \times U(1)_{B-L} \times U(1)_{I_{3R}} \times U(1)_{3B+L}$, with $U(1)_{B-L}$, $U(1)_{I_{3R}}$, and $U(1)_{3B+L} = -U(1)_X$ as defined in Eq. (3.1). Since both $U(1)_{B-L}$ and $U(1)_{3B+L}$ remain good symmetries, one immediately sees why none of the B and L-violating operators that mediate proton decay in the MSSM are allowed: Equation (1.2) is equivalent to the trivial identities

$$U(1)_B = \frac{1}{4}[U(1)_{B-L} + U(1)_{3B+L}], \qquad U(1)_L = \frac{1}{4}[-3U(1)_{B-L} + U(1)_{3B+L}], \qquad (3.3)$$

that, when acting on MSSM fields, serve the promised role of precisely counting baryon and lepton number, respectively. The relation between the global and gauged forms is given by

$$\begin{aligned} U(1)_B &= U(1)_{\text{baryon}} + U(1)_\Delta\,, \\ U(1)_L &= U(1)_{\text{lepton}} + U(1)_\Delta\,, \end{aligned} \qquad (3.4)$$

where

$$U(1)_\Delta = -\frac{1}{2}\left[3U(1)_{\text{baryon}} + U(1)_{\text{lepton}} + U(1)_{2L} + U(1)_{\text{up}} + U(1)_{\text{down}}\right] \qquad (3.5)$$

is also an anomalous $U(1)$, but under which no MSSM field is charged.

The fact that two equivalent combinations of $U(1)_{B-L}$ and $U(1)_{3B+L}$ precisely count B and L indicates that both $U(1)_B$ and $U(1)_L$ remain gauged as long as both parent symmetries remain unbroken. One may ask how the full gauge symmetry may be further broken to yield the individual combinations shown in Eq. (3.3). As an example, first note that the $U(1)_{B-L} \times U(1)_{I_{3R}}$ gauge symmetry may be broken to the SM hypercharge $U(1)_Y$ by assigning VEVs to the vectorlike singlet fields Φ, $\overline{\Phi}$ (from the $a_2 c_2'$ intersections) carrying the quantum numbers $(\mathbf{1}, \mathbf{1}, \frac{1}{2}, -1, -3)$ and $(\mathbf{1}, \mathbf{1}, -\frac{1}{2}, 1, 3)$ under the $SU(3)_C \times SU(2)_L \times U(1)_{I_{3R}} \times U(1)_{B-L} \times U(1)_{3B+L}$ gauge symmetry. Since Φ, $\overline{\Phi}$ carry nonzero charges under both $U(1)_{B-L}$ and $U(1)_{I_{3R}}$, the two gauge symmetries are broken, but the fields are uncharged under the linear combination

$$\begin{aligned} U(1)_Y &= \frac{1}{6}\left[U(1)_{\text{baryon}} - 3U(1)_{\text{lepton}} + 3U(1)_{\text{up}} - 3U(1)_{\text{down}}\right] \qquad (3.6) \\ &= \frac{1}{6}\left[U(1)_{a1} - 3U(1)_{a2} + 3U(1)_{c1} - 3U(1)_{c2}\right] \\ &= \frac{1}{2}U(1)_{B-L} + U(1)_{I_{3R}}\,, \end{aligned}$$

corresponding to the SM hypercharge, which then remains unbroken. In addition, these vectorlike fields carry charges under the $U(1)_{3B+L}$ gauge symmetry and so it is also broken.

However, these fields do not carry charges under the linear combination corresponding to $U(1)_L$ given in Eq. (3.3), and thus it survives. Therefore, at this stage, the gauge symmetry of the observable sector is given by

$$SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_L \,, \qquad (3.7)$$

just the gauge symmetry of the MSSM, extended by an extra $U(1)$ that counts L.

Many alternate scenarios for symmetry breaking are possible. For example, the $U(1)_{B-L} \times U(1)_{I_{3R}} \times U(1)_{3B+L}$ gauge symmetry may instead be broken by assigning VEVs to the right-handed neutrino fields N_R. In this case, the gauge symmetry is broken to

$$SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_B \,. \qquad (3.8)$$

However, assigning VEVs to N_R breaks SUSY, which is expected not to occur until the TeV scale. To avoid such outcomes, the $U(1)_{I_{3R}} \times U(1)_{B-L} \times U(1)_{3B+L}$ symmetry breaking is thus naively expected to occur near the usual GUT scale. Indeed, as shown in [34], if the PS gauge symmetry is broken to the MSSM at the GUT/string scale, then the MSSM gauge couplings are unified. Thus, it is natural to assume that the Φ, $\overline{\Phi}$ fields from the $a_2 c_2'$ sector obtain GUT-scale VEVs. That being said, the possibility of breaking the gauge symmetry at the TeV-scale by assigning VEVs to N_R and not Φ, $\overline{\Phi}$ is worth exploring.

IV. PHENOMENOLOGICAL CONSTRAINTS ON SINGLET VEVS

In the previous section we demonstrated that the gauge symmetry may be broken to $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_L$ at the GUT scale by assigning VEVs to the vectorlike fields Φ, $\overline{\Phi}$, or to $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_B$ at the TeV scale by assigning VEVs to the right-handed neutrinos N_R. These two cases show that models of this type may be adapted to provide either a $U(1)_L$ or $U(1)_B$ that survives unbroken to low energies.

However, other singlet fields appear in the model whose VEVs may break $U(1)_{3B+L}$ [or equivalently, $U(1)_B$ and $U(1)_L$] at intermediate scales, namely, the singlets S_L and S_R, as well as the $SU(2)_R$ triplet fields T_R. In order to determine the scales at which these fields may obtain VEVs, one must study other observable phenomena mediated by these fields. As we shall see below, one can construct operators that allow proton decay, which imposes constraints upon the VEVs of the S_R fields. In addition, one can construct superpotential

operators for the Higgsino bilinear μ term and for the Majorana mass term for right-handed neutrinos, which are in turn constrained.

A. Proton Decay via Exotic Fields

Although the $d = 4$ and $d = 5$ operators involving only MSSM fields that could mediate proton decay are excluded by the gauged $U(1)_{B-L}$ and $U(1)_{3B+L}$ gauge symmetries, the model also contains fields that are not part of the MSSM, yet are charged under $U(1)_B$ and/or $U(1)_L$. Operators involving these extra fields can mediate B- or L-violation. Furthermore, B- and L-violation can occur below energy scales at which $U(1)_{B-L}$ and $U(1)_{3B+L}$ are spontaneously broken [and, as discussed previously, if such operators are perturbatively forbidden by the global $U(1)$ symmetries, then additional instanton suppression factors appear]. One must examine the possible operators that arise and their VEVs in order to fully address these issues.

The leading such operator involving exotic fields that mediates proton decay is

$$\frac{\kappa}{M_{\mathrm{St}}^2} \, QQQLS_R \,, \tag{4.1}$$

where $M_{\mathrm{St}} = O(10^{18} \ \mathrm{GeV})$ is the string scale. Even though it is a singlet under all gauge groups, this operator is perturbatively forbidden by the global symmetries $U(1)_{4,2L,2R}$ and can only be induced via nonperturbative effects such as from D-brane instantons, which introduce the suppression factor $\kappa \propto e^{-S_{E2}}$. This operator is effectively $d = 5$, assuming that the singlet fields S_R^i obtain VEVs:

$$\frac{e^{-S_{E2}} \langle S_R^i \rangle}{M_{\mathrm{St}}^2} \, QQQL \,. \tag{4.2}$$

Using the phenomenological bound (*e.g.*, [7]) on the operator $QQQL$, we require

$$\frac{e^{-S_{E2}} \langle S_R^i \rangle}{M_{\mathrm{St}}} \leq O(10^{-7}) \,, \tag{4.3}$$

in order to satisfy the experimental lower limit on the proton lifetime. This estimate, of course, assumes the existence of a suitable D-brane instanton with the appropriate zero-mode structure in the model, which is far from obvious. Assuming such an instanton exists in the worst-case scenario $e^{-S_{E2}} \approx 1$, then the maximum allowed VEV is $\langle S_R \rangle \approx 10^{11} \ \mathrm{GeV}$; at the other extreme, an instanton suppression of 10^{-7} would allow $\langle S_R^i \rangle$ to be as large as M_{St}. A

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thorough analysis, which we do not undertake here, would require enumerating operators that, upon receiving VEVs, effectively assume the forms of the "dangerous" operators in Eqs. (2.1)–(2.3); however, the approach is completely analogous to that for Eq. (4.2).

Of course, the most straightforward way to eliminate the possibility of rapid proton decay via such operators involving exotic fields such is to forbid these fields from obtaining large VEVs. However, other phenomenological properties can depend on VEVs of these fields. Among these, in particular, are the μ term and the Majorana mass term for right-handed neutrinos, which we discuss next.

B.   The μ Term and Majorana Neutrino Masses

The problem of generating a μ superpotential term of the form $\mu H_u H_d$, with $\mu = O(\text{TeV})$, is well known in supersymmetric models. Specifically, the question of why the μ parameter is TeV-scale rather than Planck-scale has long been an open question in the MSSM. In the context of intersecting D-brane models, this problem is somewhat ameliorated in that a simple bilinear coupling is forbidden by global U(1) symmetries [here, U(1)$_{2L,2R}$]. Such a term can therefore only be generated either by high-order couplings or by nonperturbative effects such as D-brane instantons [43]. In either case, one expects the effective operator to be suppressed by powers of M_{St}.

One possibility for generating a μ term is through the higher-dimensional ($d = 5$) operator

$$W \;\supset\; \frac{y_\mu^{ijkl}}{M_{\text{St}}} S_L^i S_R^j H_u^k H_d^l \;, \tag{4.4}$$

where y_μ^{ijkl} are Yukawa couplings. In this case, the singlets S_R may obtain string or GUT-scale VEVs (or lower, as described previously) while preserving the D-flatness of U(1)$_{2R}$, and the singlets S_L may obtain TeV-scale VEVs while preserving the D-flatness of U(1)$_{2L}$, while the Higgses couple through their electroweak-scale VEVs. Simple order-of-magnitude estimates then show that a TeV-scale μ term may be generated, with $y_\mu^{ijkl} = O(1)$.

For a μ term of the desired magnitude to be generated by this higher-dimensional operator, the fields S_R must obtain VEVs near M_{St}. However, as discussed in the previous section, the constraint $\langle S_R^i \rangle \ll M_{\text{St}}$ is expected from proton decay unless strong instanton suppressions are present. Thus, while this class of models does provide a mechanism to generate μ of natural size, it is somewhat preferable for S_R to obtain VEVs at a scale much

lower than M_{St}, and for the (bulk of the) μ term to be generated via some other mechanism, such as D-brane instantons.

From Table I and Eqs. (3.1), (3.2), and (3.3), one may check that S_R^i is a singlet under $B-L$ but is charged under $3B+L$, and therefore $\langle S_R^i \rangle$ spontaneously breaks $\text{U}(1)_L$ (and similarly for $\langle T_R^i \rangle$). One then expects the scale at which S_R and T_R obtain VEVs to be correlated with the scale at which the right-handed neutrinos obtain Majorana masses, an idea that is implemented as follows. A simple Majorana mass term for right-handed neutrinos of the form $M \cdot N^c N^c$ is perturbatively forbidden in the model by the global $\text{U}(1)$ symmetries [here, $\text{U}(1)_{4,2R}$]. A Majorana mass can therefore only be generated by high-dimensional operators or by nonperturbative effects. Indeed, the possibility of generating a Majorana mass term via D-brane instantons has been much studied in the literature [46, 47]. However, it should be mentioned that for the models studied in [33, 34, 48, 49], a standard Majorana mass term $M \cdot N^c N^c$ is not a singlet under $\text{U}(1)_{B-L}$ [or, for that matter, $\text{U}(1)_{3B+L}$] and therefore cannot be generated solely via D-brane instantons. On the other hand, Majorana masses can be generated in these models without instantons by high-dimensional operators such as

$$W \supset \frac{y_{Nij}^{mnkl}}{M_{\text{St}}^3} T_R^m T_R^n \Phi^i \Phi^j F_R^k F_R^l \ , \tag{4.5}$$

where y_{Nij}^{mnkl} are Yukawa couplings, and the fields T_R and Φ may obtain VEVs at the string scale (or lower). For the model of Table I, right-handed neutrino masses can be generated in the range 10^{10-14} GeV for $y_{Nij}^{mnkl} \sim 10^{(-7)-(-3)}$, assuming GUT- or string-scale VEVs for the Φ and T_R. As noted above, $\langle T_R^i \rangle$ also break $\text{U}(1)_{3B+L}$ and $\text{U}(1)_L$, which suggests that lower values of $\langle T_R^i \rangle$ (comparable to or perhaps somewhat larger than those of $\langle S_R^i \rangle$), say, 10^{14} GeV) allow for y_{Nij}^{mnkl} to be of a natural $O(1)$ size.

In summary, if the μ term and the Majorana mass term are generated by high-dimensional operators involving the fields S_R and T_R with string-scale VEVs, then $\text{U}(1)_{3B+L}$ is broken near M_{St}. In this case, the only Abelian factor that survives at low energies is the weak hypercharge $\text{U}(1)_Y$, and the suppression of proton decay and the Majorana mass scale require small coefficients, but the μ term is of natural size. If $\langle S_R^i \rangle$ and $\langle T_R^i \rangle$ lie closer to the Majorana mass scale, then proton decay and Majorana mass terms can have natural coefficients, but then the μ term must be generated by effects other than high-dimensional operators, in particular via D-brane instanton-induced operators. Thus, it is quite possible that the μ term and Majorana mass term may be generated without requiring string-scale VEVs for

the fields S_R and T_R. If this is the case, then U(1)$_L$ can remain unbroken down to the electroweak scale.

C. Baryogenesis

The existence of B-violating processes is one of the three Sakharov conditions [50] that must be satisfied in order for baryogenesis to be realized. As shown for this model, processes that can mediate proton decay are strongly suppressed. Indeed, the relevant operators are forbidden by U(1)$_X$ and/or perturbatively forbidden by global U(1) symmetries. In fact, if global U(1)$_{\text{baryon}}$ must be conserved [*i.e.*, if no nonperturbative U(1)$_{\text{baryon}}$-violating interaction can be generated] and the nonzero VEVs are the aforementioned $\langle \Phi_{a2c2'} \rangle$, $\langle S_L \rangle$, $\langle S_R \rangle$, and $\langle T_R \rangle$ [all of which are singlets under U(1)$_{\text{baryon}}$], B-violation is forbidden. This result holds for both the single-operator and two-operator proton-decay combinations. The proton is effectively stable under these assumptions. On the other hand, since all of the VEVs listed above break global U(1)$_{\text{lepton}}$, one can show that *every* L-violating, B-conserving operator out of the 36 can be produced with the four given VEVs without violating any symmetry. Thus, L-violating processes clearly dominate over B-violating ones in this model.

The most likely possibility for realizing the matter-antimatter symmetry of the universe within this model is therefore electroweak baryogenesis [51, 52]. In this scenario, leptogenesis occurs first, thereby generating a lepton-number asymmetry. Nonperturbative field configurations known as sphalerons then convert this lepton-number asymmetry into a baryon-number asymmetry, resulting in baryogenesis. Our model allows numerous L-violating operators at the perturbative level, while B-violation requires instanton-like configurations to occur; even if such configurations are heavily suppressed in the modern universe, they may be much easier to realize in the high-temperature environment of the early universe. Electroweak baryogenesis is therefore touted as a natural means of generating a baryon asymmetry while also maintaining the effective stability of the proton.

One should recall that a sufficient amount of CP violation (CPV) is also a necessary Sakharov condition for baryogenesis. While the minimal three-generation SM is well known to possess insufficient CPV to account for baryogenesis, the introduction of a fourth generation [53] provides a suitable natural source of CPV. Other CPV sources enter through the MSSM two-Higgs potential, and indeed, through the couplings of numerous exotic fields

appearing in our string-inspired model. All of the basic ingredients necessary to realize the observed baryogenesis therefore appear to be present in this model.

V. PROSPECTS FOR NEW PHYSICS

The four-generation model (or more accurately, family of models) discussed here was constructed in Type IIA string theory with D6-branes intersecting at angles. As such, the model satisfies all global consistency constraints as well as preserving $\mathcal{N} = 1$ supersymmetry. In addition, the model has a rich phenomenology, which includes rank-2 Yukawa matrices that can provide a natural explanation for the mass textures required to satisfy current constraints on a fourth generation of chiral fermions. Also, the MSSM gauge couplings are unified at the tree-level in the model, and matter charged under the hidden sector decouples at high energies. These features by themselves make this string theory model of high phenomenological interest, and provide additional motivation for the possible existence of a fourth generation of quarks and leptons that might be observed at the LHC.

Given the above features, it is really quite remarkable that the model automatically possesses an extra nonanomalous U(1) gauge symmetry that forbids all of the dimension-4, -5, and -6 operators that mediate proton decay in the MSSM. Although the idea of eliminating these operators using an extra nonanomalous U(1) is not new, the extra U(1) of this model is not added to the model by hand, but rather emerges automatically from the D-brane construction. Specifically, this extra U(1) has not been made anomaly-free by arbitrarily adding extra matter representations into the model as is done in most phenomenological models in the literature. Among the phenomenological consequences of the existence of this U(1) is that the proton is, for all intents and purposes, stable in the model. This result is consistent with the non-observation of proton decay, and furthermore implies that proton decay should not be observed in future experiments. This being said, as discussed in the previous section, proton decay may be allowed in principle via operators involving exotic matter fields. However, these operators are forbidden by global symmetries in the model, and it has not been demonstrated that they may be induced by nonperturbative effects such as D-brane instantons. Even if such effects occur, one then merely obtains constraints on the VEVs of the exotic singlet fields in the model, as discussed above.

As the dimension-4 operators that mediate rapid proton decay are forbidden by the

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Abelian gauge symmetries, the role for an R parity is much diminished in the model; R parity is not required to eliminate rapid proton decay in this model as it does in the MSSM. This fact may have dramatic consequences for the supersymmetric phenomenology at the LHC and for dark matter experiments. In particular, the absence of R parity results in an unstable Lightest Supersymmetric Particle (LSP). In the MSSM with R parity, the LSP is stable and thus must be electromagnetically neutral, which provides a natural candidate for dark matter. In addition, many of the collider signatures that provide the cleanest signals for the production of supersymmetric particles are those with large missing transverse energy (MET). Such signals are expected to arise when the produced supersymmetric particles decay directly to the LSP, which is typically a neutralino. However, if the LSP is no longer stable in the absence of R parity, then these MET signals do not have the same significance. Indeed, in this scenario the LSP is not even required to be electromagnetically neutral. The lack of R parity makes the observation of supersymmetric particles at the LHC a much more difficult proposition. In addition, the LSP cannot be stable dark matter, which would strongly impact the prospects for direct dark matter detection experiments.

Finally, as we have seen, it is possible for the gauge symmetry to be broken to the MSSM with an extra gauged U(1) that couples either to lepton [U(1)$_L$] or to baryon [U(1)$_B$] number. Whether one obtains a gauged lepton number or a gauged baryon number at low energies depends intrinsically upon which fields are used to break the gauge symmetry. Either U(1)$_L$ or U(1)$_B$ may survive unbroken to the TeV scale or below. If so, the symmetry would be heralded by so-called leptophilic or leptophobic Z' bosons, respectively, which would be observable at the LHC. As mentioned above, a leptophilic Z' could explain [12] the observed PAMELA [13]/ATIC [14]/Fermi LAT [15] cosmic ray positron excess, while a leptophobic Z' has been discussed in connection with recent anomalies observed by the CDF collaboration [16–18]. Such considerations deserve detailed study, which we leave for future work.

VI. CONCLUSIONS

We have demonstrated the existence of an extra massless and anomaly-free U(1) gauge symmetry in the model constructed in [34] that plays a central role in suppressing B- and L-violation. In particular, we obtained linear combinations of this U(1)$_X$ with U(1)$_{B-L}$ that

count baryon and lepton number, respectively, thereby providing $U(1)_{B,L}$ gauge symmetries at the string scale that forbid all dimension-4, dimension-5, and dimension-6 operators involving only MSSM fields that mediate proton decay.

Nevertheless, additional fields charged under the gauged B and L symmetries appear in the model; other operators involving this exotic matter arise that can mediate proton decay through their VEVs, once $U(1)_X$ breaks spontaneously. We found the leading such operator $QQQLS_R$, and while noting that it can only appear nonperturbatively in the model, estimated the constraint on the operator's coefficient. Constrained, in turn, is the scale at which the singlet fields S_R may receive VEVs. We noted that VEVs for S_R break $U(1)_L$; if $\langle S_R \rangle$ is low, as suggested by the proton decay constraint, then $U(1)_L$ can potentially remain unbroken down to the electroweak scale. If so, this scenario can provide a so-called leptophilic Z' boson, which could be observed at the LHC and which can explain the galactic cosmic ray positron excess observed by the PAMELA collaboration.

The appearance of this extra massless and anomaly-free gauge symmetry $U(1)_X$ in the model is quite remarkable. This extra $U(1)$ occurs naturally alongside the usual MSSM gauge symmetry, and is therefore arguably a natural extension of the MSSM. The fact that $U(1)_X$ forbids all of the well-known dimension-4 and -5 operators that can mediate rapid proton decay, while also allowing all of the Yukawa couplings for quarks and leptons, is especially remarkable. As we have seen, this extra Abelian gauge symmetry can be viewed as a means of gauging baryon and lepton number, which is strongly compatible with the non-observation of B- and L-violating effects such as proton decay and neutrinoless double beta decay. These features, combined with previous results shown for this particular model that the MSSM gauge couplings unify at the string scale and that it is possible to naturally obtain realistic Yukawa mass matrices for quarks and leptons, give this model particular phenomenological value.

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