PHYSICAL REVIEW D **84,** 106009 (2011)

Proton stability and dark matter in a realistic string MSSM

James A. Maxin,[1] Van E. Mayes,[2] and D. V. Nanopoulos[1,3]

[1]*George P. and Cynthia W. Mitchell Institute for Fundamental Physics, Texas A&M University, College Station, Texas 77843, USA*
[2]*Physics Department, Arizona State University, Tempe, Arizona 85287-4111, USA*
[3]*Astroparticle Physics Group, Houston Advanced Research Center (HARC), Mitchell Campus, Woodlands, Texas 77381, USA;*
Academy of Athens, Division of Natural Sciences, 28 Panepistimiou Avenue, Athens 10679, Greece

(Received 21 August 2011; published 11 November 2011)

We demonstrate the existence of an extra nonanomalous $U(1)$ gauge symmetry in a three-generation Pati-Salam model constructed with intersecting D6-branes in Type-IIA string theory on a $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold. This extra $U(1)$ forbids all dimension-4, 5, and 6 operators which mediate proton decay in the MSSM. Moreover, this results in the effective promotion of baryon and lepton number to local gauge symmetries, which can potentially result in leptophobic and leptophilic Z' bosons observable at the LHC. Furthermore, it is not necessary to invoke R-parity to forbid the dimension-4 operators which allow rapid proton decay. However, R-parity may arise naturally from a spontaneously broken $U(1)_{B-L}$. Assuming the presence of R-parity, we then study the direct-detection cross-sections for neutralino dark matter, including the latest constraints from the XENON100 experiment. We find that these limits are now within required range necessary to begin testing the model.

DOI: 10.1103/PhysRevD.84.106009 PACS numbers: 11.25.Mj

I. INTRODUCTION

A main goal of string phenomenology is to discover the mechanisms by which the detailed properties of our universe may arise. Among these are the standard model (SM) gauge groups, the number of generations of chiral fermions, and the observed mass hierarchies and mixings of quarks and leptons. Of particular importance, the proton appears to have a very long lifetime. Baryon (B) and lepton number (L) violating processes have to date never been observed, yet they are only conserved as accidental global symmetries of the SM. However, such global symmetries are generically broken by nonperturbative effects, and thus baryon number is expected to be violated at some level in the SM. In fact, dimension-4 operators appear in the minimal supersymmetric SM (MSSM) which leads to proton decay at a disastrously high rate. Usually these operators are eliminating by imposing a discrete symmetry on the MSSM such as R-parity [1–3]. However, operators of dimension-5 also appear which are not eliminated by R-parity, leading to a proton decay rate which is generically too large.

Although imposition of R-parity may seem *ad hoc*, it provides a simple explanation for another mystery. Observations in cosmology and astrophysics suggest the presence of a stable dark matter particle. A natural candidate for WIMP-like dark matter is the lightest supersymmetric partner (LSP) [4] in supersymmetric models which include R-parity conservation, which is usually the lightest neutralino $\tilde{\chi}_1^0$ [4,5]. Limits on the dark matter relic abundance and direct and indirect detection cross sections can be used to constrain the possible superpartner and Higgs spectra, which may be observed at the Large Hadron Collider (LHC). In short, only superpartner spectra which possess a stable LSP consistent with all other constraints on dark matter are viable.

Of course, the MSSM is just an effective theory which should be replaced at high energies by something more fundamental, such as string theory. In particular, Type-IIA string compactifications involving D6-branes intersecting at angles (and their Type IIB duals including F-theory extensions) have provided a fruitful direction for studying this question. Such models have been the subject of much study in recent years, and we refer the reader to [6,7] for recent reviews. A phenomenologically interesting model of this type was first constructed in [8,9] and studied in [9–11]. In this three-generation Pati-Salam model, it is possible to obtain realistic Yukawa matrices for quarks and leptons, tree-level gauge unification at the string-scale, and obtain realistic supersymmetry spectra satisfying all experimental constraints. The phenomenological consequences of this model at the LHC were considered in [11,12], and the implications for direct and indirect dark matter detection were initially studied in [13,14]. In the present work, we show that a variation of this model originally constructed in [15] possesses an extra nonanomalous $U(1)$ gauge symmetry that forbids all dimension-4, 5, and 6 operators found in the MSSM which allow proton decay (related four-generation models were considered in [16–18]). Thus, the proton is effectively stable in the model (see [19] for a similar study in the context of free-fermionic heterotic string/M-theory compactifications). In particular, it is not necessary to introduce R-parity in order to eliminate the dimension-4 operators which allow proton decay at a dangerously high rate. Nevertheless, R-parity may still naturally arise in the model via a $U(1)_{B-L}$ gauge symmetry, which is broken spontaneously to its discrete $\mathbb{Z}_2$ subgroup, resulting in a stable LSP. Thus, we update the

TABLE I. General spectrum for intersecting D6 branes at generic angles, where $I_{aa'} = -2^{3-k}\prod_{i=1}^{3}(n_a^i l_a^i)$ and $I_{aO6} = 2^{3-k}(-l_a^1 l_a^2 l_a^3 + l_a^1 n_a^2 n_a^3 + n_a^1 l_a^2 n_a^3 + n_a^1 n_a^2 l_a^3)$, where $k = \beta_1 + \beta_2 + \beta_3$. In addition, $\mathcal{M}$ is the multiplicity, and a_S and a_A denote the symmetric and antisymmetric representations of $U(N_a/2)$, respectively.

Sector	Representation
aa	$U(N_a/2)$ vector multiplet and 3 adjoint chiral multiplets
$ab + ba$	$\mathcal{M}(\frac{N_a}{2}, \frac{\bar{N}_b}{2}) = I_{ab} = 2^{-k}\prod_{i=1}^{3}(n_a^i l_b^i - n_b^i l_a^i)$
$ab' + b'a$	$\mathcal{M}(\frac{N_a}{2}, \frac{N_b}{2}) = I_{ab'} = -2^{-k}\prod_{i=1}^{3}(n_a^i l_b^i + n_b^i l_a^i)$
$aa' + a'a$	$\mathcal{M}(a_S) = \frac{1}{2}(I_{aa'} - \frac{1}{2}I_{aO6})$; $\mathcal{M}(a_A) = \frac{1}{2}(I_{aa'} + \frac{1}{2}I_{aO6})$

constraints on dark matter direct-detection taking into account the recent limit on the dark matter direct-detection cross-section from the CDMSII [20] and XENON100 [21] collaborations.

II. A REALISTIC MSSM WITH AN EXTRA $U(1)$

Type IIA orientifold string compactifications with intersecting D-branes (and their Type IIB duals with magnetized D-branes) have provided exciting geometric tools with which the MSSM may be engineered. While this approach may not allow a first-principles understanding of why the SM gauge groups and associated matter content arises, it may allow a deeper insight into how the finer phenomenological details of the SM may emerge. In short, D6-branes in Type IIA fill (3 + 1)-dimensional Minkowski spacetime and wrap 3-cycles in the compactified manifold, such that a stack of N branes generates a gauge group $U(N)$ [or $U(N/2)$ in the case of $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$] in its world volume. On $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$, the 3-cycles are of the form [22]

$$\Pi_a = \prod_{i=1}^{3}(n_a^i[a_i] + 2^{-\beta_i} l_a^i[b_i]), \tag{1}$$

where the integers n_a^i and l_a^i are the wrapping numbers around the basis cycles $[a_i]$ and $[b_i]$ of the ith two-torus, and $\beta_i = 0$ for an untilted two-torus, while $\beta_i = 1$ for a tilted two-torus. In addition, we must introduce the orientifold images of each D6-brane, which wraps a cycle given by

$$\Pi_a' = \prod_{i=1}^{3}(n_a^i[a_i] - 2^{-\beta_i} l_a^i[b_i]). \tag{2}$$

In general, the 3-cycles wrapped by the stacks of D6-branes intersect multiple times in the internal space, resulting in a chiral fermion in the bifundamental representation localized at the intersection between different stacks a and b. The multiplicity of such fermions is then given by the number of times the 3-cycles intersect. Each stack of D6-branes a may intersect the orientifold images of other stacks b', also resulting in fermions in bifundamental representations. Each stack may also intersect its own image a', resulting in chiral fermions in the symmetric and antisymmetric representations. The different types of representations that may be obtained for each type of intersection and their multiplicities are summarized in Table I. In addition, the consistency of the model requires certain constraints to be satisfied, namely, Ramond-Ramond (R-R) tadpole cancellation and the preservation of $\mathcal{N} = 1$ supersymmetry.

The set of D6 branes wrapping the cycles on a $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold shown in Table II results in a three-generation Pati-Salam model with additional hidden sectors. The full gauge symmetry of the model is given by $[U(4)_C \times U(2)_L \times U(2)_R]_{\text{observable}} \times [U(2) \times USp(2)^2]_{\text{hidden}}$, with the matter content shown in Table III. As discussed in detail in [10,11], with this configuration of D6 branes all R-R tadpoles are canceled, K-theory constraints are satisfied, and $\mathcal{N} = 1$ supersymmetry is preserved. Furthermore, the tree-level MSSM

TABLE II. D6-brane configurations and intersection numbers for a three-family Pati-Salam model on a Type-IIA $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold, with a tilted third two-torus. The complete gauge symmetry is $[U(4)_C \times U(2)_L \times U(2)_R]_{\text{observable}} \times [U(2) \times USp(2)^2]_{\text{hidden}}$ and $\mathcal{N} = 1$ supersymmetry is preserved for $\chi_1 = 3$, $\chi_2 = 1$, $\chi_3 = 2$.

	$U(4)_C \times U(2)_L \times U(2)_R \times U(2) \times USp(2)^2$											
	N	$(n^1, l^1) \times (n^2, l^2) \times (n^3, l^3)$	n_S	n_A	b	b'	c	c'	d	d'	3	4
a	8	$(0, -1) \times (1, 1) \times (1, 1)$	0	0	3	0	-3	0	0(2)	0(1)	0	0
b	4	$(3, 1) \times (1, 0) \times (1, -1)$	2	-2	-	-	0(6)	0(1)	1	0(1)	0	-3
c	4	$(3, -1) \times (0, 1) \times (1, -1)$	-2	2	-	-	-	-	-1	0(1)	3	0
d	4	$(1, 0) \times (1, -1) \times (1, 1)$	0	0	-	-	-	-	-	-	-1	1
3	2	$(0, -1) \times (1, 0) \times (0, 2)$					$\chi_1 = 3$					
4	2	$(0, -1) \times (0, 1) \times (2, 0)$					$\chi_2 = 1$, $\chi_3 = 2$					

TABLE III. The chiral and vectorlike superfields, their multiplicities and quantum numbers under the gauge symmetry $[U(4)_C \times U(2)_L \times U(2)_R]_{\text{observable}} \times [U(2) \times USp(2)^2]_{\text{hidden}}$, where $Q_X = Q_4 + 2(Q_{2L} + Q_{2R} + 3Q_d)$.

	Mult.	Quantum Number	Q_4	Q_{2L}	Q_{2R}	Q_X	Field
ab	3	$(4, \bar{2}, 1, 1, 1, 1)$	1	−1	0	−1	$F_L(Q_L, L_L)$
ac	3	$(\bar{4}, 1, 2, 1, 1, 1)$	−1	0	1	1	$F_R(Q_R, L_R)$
bd	1	$(1, \bar{2}, 1, 2, 1, 1)$	0	−1	0	4	X_{bd}
cd	1	$(1, 1, 2, \bar{2}, 1, 1)$	0	0	1	−4	X_{cd}
$b4$	3	$(1, \bar{2}, 1, 1, 1, 2)$	0	−1	0	−2	X_{b3}^i
$c3$	3	$(1, 1, 2, 1, \bar{2}, 1)$	0	0	1	2	X_{c3}^i
$d3$	1	$(1, 1, 1, \bar{2}, 2, 1)$	0	0	1	−6	X_{cd}
$d4$	1	$(1, 1, 1, 2, 1, \bar{2})$	0	0	1	6	X_{cd}
b_S	2	$(1, 3, 1, 1, 1, 1)$	0	2	0	4	T_L^i
b_A	2	$(1, \bar{1}, 1, 1, 1, 1)$	0	−2	0	−4	S_L^i
c_S	2	$(1, 1, \bar{3}, 1, 1, 1)$	0	0	−2	−4	T_R^i
c_A	2	$(1, 1, 1, 1, 1, 1)$	0	0	2	4	S_R^i
ab'	3	$(4, 2, 1, 1, 1, 1)$	1	1	0	3	Ω_L^i
	3	$(\bar{4}, \bar{2}, 1, 1, 1, 1)$	−1	−1	0	−3	$\bar{\Omega}_L^i$
ac'	3	$(4, 1, 2, 1, 1, 1)$	1	0	1	3	Φ_i
	3	$(\bar{4}, 1, \bar{2}, 1, 1, 1)$	−1	0	−1	−3	$\bar{\Phi}_i$
bc	6	$(1, 2, \bar{2}, 1, 1, 1)$	0	1	−1	0	H_u^i, H_d^i
	6	$(1, \bar{2}, 2, 1, 1, 1)$	0	−1	1	0	

gauge couplings are unified at the string scale. Finally, the Yukawa matrices for quarks and leptons are rank 3 and it is possible to obtain correct mass hierarchies and mixings. Note that the observable sector of the model shown in Tables II and III is identical to that of references [10,11] so that all of the above phenomenological features are also present. However, the hidden sector of the model is different, which as we shall see gives rise to an extra anomaly-free $U(1)$ gauge symmetry.

Since $U(N) = SU(N) \times U(1)$, associated with each the stacks a, b, c, and d are $U(1)$ gauge groups, denoted as $U(1)_a$, $U(1)_b$, $U(1)_c$, and $U(1)_d$. In general, these $U(1)$s are anomalous. The anomalies associated with these $U(1)$s are canceled by a generalized Green-Schwarz (G-S) mechanism that involves untwisted R-R forms. The couplings of the four untwisted R-R forms B_2^i to the $U(1)$ field strength F_a of each stack a are given by [22,23]

$$\begin{aligned} &N_a l_a^1 n_a^2 n_a^3 \int_{M4} B_2^1 \wedge \mathrm{tr} F_a, \quad N_a n_a^1 l_a^2 n_a^3 \int_{M4} B_2^2 \wedge \mathrm{tr} F_a, \\ &N_a n_a^1 n_a^2 l_a^3 \int_{M4} B_2^3 \wedge \mathrm{tr} F_a, \quad -N_a l_a^1 l_a^2 l_a^3 \int_{M4} B_2^4 \wedge \mathrm{tr} F_a. \end{aligned} \tag{3}$$

As a result, the gauge bosons of these Abelian groups generically become massive. However, these $U(1)$s remain as global symmetries to all orders in perturbation theory. Indeed, baryon and lepton number conservation are typically identified as arising from these global symmetries. These global $U(1)$ symmetries may also result in the forbidding of certain superpotential operators, such as Yukawa couplings and those which mediate baryon and lepton number violation. However, these *global* symmetries may be broken by nonperturbative effects, such as from D-brane instantons.

The couplings of Eq. (3) determine the exact linear combinations of $U(1)$ gauge bosons that acquire string-scale masses via the G-S mechanism. If $U(1)_X$ is a linear combination of the $U(1)$s from each stack,

$$U(1)_X \equiv \sum_a C_a U(1)_a, \tag{4}$$

then the corresponding field strength must be orthogonal to those that acquire G-S mass. Thus, if a linear combination $U(1)_X$ satisfies [23–25]

$$\begin{aligned} &\sum_a C_a N_a l_a^1 n_a^2 n_a^3 = 0, \qquad \sum_a C_a N_a n_a^1 l_a^2 n_a^3 = 0, \\ &\sum_a C_a N_a n_a^1 n_a^2 l_a^3 = 0, \qquad \sum_a C_a N_a l_a^1 l_a^2 l_a^3 = 0, \end{aligned} \tag{5}$$

the gauge boson of $U(1)_X$ acquires no G-S mass and is anomaly-free, provided that the RR-tadpole conditions are satisfied.

For the present model, precisely one linear combination satisfies the above conditions, and therefore has a massless gauge boson and is anomaly- free:

$$U(1)_X = U(1)_a + 2[U(1)_b + U(1)_c + 3U(1)_d]. \tag{6}$$

Thus, the effective gauge symmetry of the model at the string scale is given by

$$\begin{aligned} SU(4)_C \times SU(2)_L \times SU(2)_R \times U(1)_X \\ \times [SU(2) \times USp(2)^2]. \end{aligned} \tag{7}$$

As can be seen from Table III, the superfields $F_L^i(Q_L, L_L)$ carry charge $Q_X = -1$, the superfields $F_R^i(Q_R, L_R)$ carry charge $Q_X = +1$, while the Higgs superfields are uncharged under $U(1)_X$. Thus, the trilinear Yukawa couplings for quarks and leptons are allowed by both the global $U(1)$ symmetries as well as the gauged $U(1)_X$ symmetry. As was shown in [10,11], the resulting Yukawa matrices are rank 3, which allows for fermion mass textures that can easily accommodate the observed mass hierarchies and mixings for quarks and leptons.

The Pati-Salam gauge symmetry is broken to the SM in two steps. First, the a and c stacks of D6-branes are split such that $a \rightarrow a1 + a2$ and $c \rightarrow c1 + c2$, where $N_{a1} = 6$, $N_{a2} = 2$, $N_{c1} = 2$, and $N_{c2} = 2$. The process of breaking the gauge symmetry via brane splitting corresponds to assigning VEVs along flat directions to adjoint scalars associated with each stack that arise from the open-string moduli [8]. After splitting the D6-branes, the gauge symmetry of the observable sector is

$$SU(3)_C \times SU(2)_L \times U(1)_{I3R} \times U(1)_{B-L} \times U(1)_{3B+L}, \tag{8}$$

where

$$\begin{aligned} U(1)_{I3R} &= \frac{1}{2}(U(1)_{c1} - U(1)_{c2}), \\ U(1)_{B-L} &= \frac{1}{3}(U(1)_{a1} - 3U(1)_{a2}), \end{aligned} \tag{9}$$

and

$$\begin{aligned} U(1)_{3B+L} = -[&U(1)_{a1} + U(1)_{a2} + 2(U(1)_b + U(1)_{c1} \\ &+ U(1)_{c2} + 3U(1)_d)], \end{aligned} \tag{10}$$

and $U(1)_{3B+L} = -U(1)_X$. Just as was the case in [18], one may also form linear combinations of $U(1)_{B-L}$ and $U(1)_{3B+L}$, which couple to baryon number and lepton number, respectively:

$$\begin{aligned} U(1)_B &= \frac{1}{4}[U(1)_{B-L} + U(1)_{3B+L}], \\ U(1)_L &= \frac{1}{4}[-3U(1)_{B-L} + U(1)_{3B+L}]. \end{aligned} \tag{11}$$

As mentioned in the Introduction, the promotion of the SM to the MSSM introduced operators which allow proton decay. The first of these is the rapid decay of the proton through the pair of $d = 4$ F-term operators (B- and L-violating, respectively) [26]:

$$U^c D^c D^c, \qquad QD^c L. \tag{12}$$

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 [27–29]. For the present model, none of these operators are singlets under either $U(1)_{B-L}$ or $U(1)_{3B+L}$ and so are forbidden.

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 $B - L$:

$$[QQQL]_F, \qquad [U^c U^c D^c E^c]_F, \qquad [D^c D^c U^c N^c]_F. \tag{13}$$

The second set consists of relevant $d = 5$ operators that violate either B or L separately, which combine with the appropriate member of Eq. (12) to form a composite operator that conserves $B - L$ and allows proton decay:

$$\begin{aligned} &[QQQH_d]_F, && [QU^c E^c H_d]_F, \\ &[QU^c L^\dagger]_D, && [U^c (D^c)^\dagger E^c]_D, \\ &[QQ(D^c)^\dagger]_D, && [QQ^\dagger N^c]_D, \\ &[U^c (U^c)^\dagger N^c]_D, && [D^c (D^c)^\dagger N^c]_D, \\ &[QU^c N^c H_u]_F, && [QD^c N^c H_d]_F. \end{aligned} \tag{14}$$

Indeed, these $d = 5$ operators are those which effectively lead to the exclusion of GUTs based on minimal $SU(5)$ [30], although these operators can be suppressed in other unified models, in particular, flipped $SU(5)$ [31–33]. For the present model, it should be noted that these operators are invariant under $U(1)_{B-L}$; however they are not invariant under $U(1)_{3B+L}$. Thus, these operators are also forbidden in the model. Similar considerations apply to the dimension-6 proton decay operators. Of course, these results may be easily understood by considering that baryon and lepton number are effectively gauged in the model as given by Eq. (11). It should also be emphasized that since these operators are forbidden by gauged symmetries rather than global symmetries, none of these operators may appear either perturbatively or nonperturbatively. Thus, the proton is essentially stable in this model with a lifetime in excess of the current experimental lower bounds.

Of course, the gauge symmetry must be further broken to the SM, with the possibility of one or more additional $U(1)$ gauge symmetries. This may be accomplished in this model by assigning VEVs to the vectorlike singlet fields with 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 from the $a_2 c_2'$ intersections. In this case, the gauge symmetry is further broken to

$$[SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_L]_{\text{observable}} \times [SU(2) \times USp(2)^2]_{\text{hidden}}, \quad (15)$$

where $U(1)_L$ is given in Eq. (11) and the electroweak hypercharge is given by the combination

$$U(1)_Y = \frac{1}{6}[U(1)_{a1} - 3U(1)_{a2} + 3U(1)_{c1} - 3U(1)_{c2}] = \frac{1}{2}U(1)_{B-L} + U(1)_{I3R}. \quad (16)$$

As we can see, if the gauge symmetry is broken to the SM in this way, $U(1)_L$ survives.

On the other hand, other 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

$$[SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_B]_{\text{observable}} \times [SU(2) \times USp(2)^2]_{\text{hidden}}. \quad (17)$$

However, assigning VEVs to N_R breaks SUSY, which is expected not to occur until the TeV scale. Thus, it is possible to obtain a nonanomalous gauged $U(1)$ which counts either lepton number or baryon number, depending upon the way in which singlet VEVs are assigned.

III. LEPTOPHOBIC AND LEPTOPHILIC Z' BOSONS

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 vector-like fields Φ, $\bar{\Phi}$, 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. The possibility of constructing models where baryon and lepton number are gauged at low energies has, of course, been considered before [34–37]. Usually in such models, extra matter must be arbitrarily added in order to cancel anomalies. For the present construction, the matter content and anomaly cancellation is fixed by the configuration of D-branes and the global consistency conditions. Thus, it is possible to obtain nonanomalous $U(1)$ gauge symmetries coupled to baryon and lepton number in a very natural way (see [38] for a discussion of the implications of extra Abelian gauge symmetries in string models).

In addition to those fields discussed above, 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 particular, the μ-term and a Majorana mass term may be generated by superpotential operators of the form

$$W \supset \frac{y_\mu^{ijkl}}{M_{\text{St}}} S_L^i S_R^j H_u^k H_d^l + \frac{y_{Nij}^{mnkl}}{M_{\text{St}}^3} T_R^m T_R^n \Phi^i \Phi^j F_R^k F_R^l, \quad (18)$$

where y_μ^{ijkl} and y_{Nij}^{mnkl} are Yukawa couplings. In this case, the singlets S_R and T_R may obtain string or GUT-scale VEVs (or lower) 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 by these operators, with $y_\mu^{ijkl} = O(1)$ and 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. In this is the case, the only surviving Abelian symmetry in the model which survives is the SM hypercharge, $U(1)_Y$. However, it is also possible to generate a μ-term and a right-handed Majorana mass via nonperturbative effects such as D-brane instantons [39]. In this case, the singlet fields S_R or T_R need not receive VEVs at high energy scales, and so either $U(1)_B$ or $U(1)_L$ may potentially survive unbroken.

If either $U(1)_B$ or $U(1)_L$ survives unbroken down to the TeV-scale, this may result in so-called leptophobic (coupled to quarks, but not leptons) or leptophilic (coupled to leptons, but not quarks) Z' bosons which may be observable at the LHC. In particular, leptophobic Z' bosons have been obtained in unified models based on flipped $SU(5)$ and E_6, though typically with couplings which are family nonuniversal [40,41]. The possibility of observing Z' bosons in general has been much studied in the literature and we direct the reader to [42,43] for reviews.

The main constraints on Z' bosons with electweak scale couplings come from precision electroweak data, direct searches at the Tevatron, and searches for flavor-changing neutral currents (FCNC). Perhaps the most stringent constraints on Z' couplings comes from LEP II. For example, the process $e^+e^- \to Z' \to e^+e^-$ leads to a constraint of $g_{eeZ'} \lesssim 0.044 \times (m_{Z'}/200$ GeV for Z' masses above roughly 200 GeV [44–46]. At lower mass scales, the LEP II constraint, which is derived in an effective field theory formalism, is not directly applicable. Below about a scale of 200 GeV, off-shell Z' production is not suppressed by the Z' mass, but instead by the LEP center-of-mass energy. A modest constraint is therefore $g_{eeZ'} \lesssim 0.04$ for $m_{Z'} \lesssim 200$ GeV. Constraints which are somewhat strong may be placed on the production and decay into e^+e^- pairs of on-shell Z' bosons if the Z' mass is near one of the center-of-mass energies at which LEP II operated [46]. Constraints from the s-channel production of e^+e^- [47] and/or $\mu^+\mu^-$ [48] at the Tevatron are also quite stringent ($\tau^+\tau^-$ final states are considerably less constrained [49]). A Z' with standard model-like couplings, for example, must be heavier than approximately 1 TeV to be consistent with the null results of these searches [50].

A so-called leptophobic Z', such as would result from $U(1)_B$, is much more difficult to observe at both lepton and hadron colliders. In particular, at hadron colliders the QCD background at low dijet mass introduces large theoretical uncertainties, overwhelming any resonance signal arising from a Z' with electroweak-strength or smaller couplings, thus the naive expectation that a search for a peak in the dijet invariant mass distributions would suffice is not correct. For a leptophobic Z' in the mass range $\sim$300–900 GeV, dijet searches at the Tevatron ($p\bar{p} \rightarrow Z' \rightarrow q\bar{q}$) constrain its couplings to quarks to be comparable to or less than those of the standard model Z [51]. For a leptophobic Z' below 300 GeV, the uncertainties in the QCD background overwhelm the signal at the Tevatron, and so the strongest constraints come from the lower energy UA2 experiment [52]. From the lack of an observed dijet resonance, UA2 can place constraints on the order of $g_{qqZ'} \lesssim 0.2$–0.5 for Z' masses in the range of 130 to 300 GeV. From these constraints, we can see that a leptophilic Z' resulting from $U(1)_L$ would require a mass greater than 1 TeV, while a leptophobic Z' resulting from $U(1)_B$ may be light so long as its couplings to quarks are comparable to the Z boson of the SM.

Both leptophilic and leptophobic Z' bosons have been put forward as explanations of various experimental anomalies in recent years. The possibility of leptophilic dark matter, such as might arise in the present context if the gaugino associated with $U(1)_L$ is stable, has been suggested as an explanation [53] of the observed PAMELA [54]/ATIC [55] cosmic ray positron excess, while a relatively light leptophobic Z' has been suggested as an explanation [56] of the Tevatron anomalies in the measured $t\bar{t}$ forward-backward asymmetry [57] and the associated production of Ws with jets [58], although a more recent analysis by the CMS collaboration has ruled out a Z' as an explanation of the forward-backward asymmetry [59]. Furthermore, the D0 collaboration has not observed the same W + dijet excess as CDF [60]. The goal of the present work is not to provide a solution for these issues, but rather to demonstrate that such Z' bosons may exist in the model and suggest possible applications (see [61,62] for a similar recent discussion in the context of Type II string compactifications with a low string scale).

As the possibility of low-scale Z' bosons has been extensively studied in the literature, including leptophilic and leptophobic varieties, it is not necessary to repeat these analyses in the present context. We have shown that the model may allow for such Z' bosons to be present at low-energies, and the results of previous studies on Z' bosons are applicable to these results. Most importantly, the Z' couplings in this model are family universal, thus they do not give rise to flavor-changing neutral currents (FCNC). Perhaps the most exciting possibility for new physics involving Z' bosons at the moment is that a leptophobiz Z' can explain the W + dijet excess reported by CDF [56,63]. Finally, let us recall that the gauge symmetry of the may be broken so that $U(1)_B$ survives below the TeV scale by assigning VEVs to the right-handed neutrino fields, N_R. As the VEVs of these fields break supersymmetry, the scale at which this is expected is the TeV scale. Thus, a leptophobic Z' in the model is expected to have mass of the TeV scale or lower, while a leptophilic Z' may have a mass intermediate between the TeV scale and the unification scale.

IV. *R*-PARITY AND NEUTRALINO DARK MATTER

As discussed in the previous section, all dimension-4, 5, and 6 operators which arise in the MSSM that may mediate proton decay are forbidden in this model by the extra $U(1)$ gauge symmetry. In the conventional MSSM, the dimension-4 operators which lead to proton decay at a disastrously high rate are typically removed by invoking R-parity. As a bonus, this results in a stable LSP, which can provide an excellent dark matter candidate in the case of a neutralino or gravitino LSP. However, as we have seen, it is not necessary to invoke R-parity in this model in order to eliminate rapid proton decay. Thus, it is possible that the LSP may not be stable in this model, and therefore would not provide a dark matter candidate. However, the model does possess a gauged $U(1)_{B-L}$ at the string scale after the Pati-Salam gauge symmetry is broken. This can then provide a natural origin for a gauged R-parity, even though it is not required for proton stability. In particular, the model will possess an exact gauged R-parity provided that $U(1)_{B-L}$ is broken by scalar VEVs that carry even integer values of $3(B-L)$ [29]. As this can clearly be accomplished in the model, in the following we will consider that an exact gauged R-parity does exist, and a stable LSP provides the required dark matter candidate. We then will analyze the constraints on possible cross sections for direct dark matter detection in light of the latest experimental data from the XENON100 and CDMS experiments. However, it should be noted that the case without exact R-parity would also be very interesting to study. In particular, this could result in different decay cascades as well as the absence of large missing energy signals since the LSP would not be stable. Needless to say, such a scenario could make it somewhat more difficult to observe superpartners at the LHC. For this reason, as well as others, it is therefore very important to also study dark matter direct-detection experiments in order to compare the predictions of supersymmetric models with the actual properties of the dark matter.

In contrast to phenomenological frameworks such as mSUGRA, the supersymmetry-breaking soft terms in intersecting D6-branes are in general nonuniversal [64]. Thus, it is possible to obtain a parameter space which is more general than in mSUGRA. A detailed discussion of the supersymmetry parameter space of the D6 model may be found in [10–12]. A comparison of the superpartner parameter space for the present case with nonuniversal soft-terms to one-parameter models motivated by no-scale supergravity may be found in [65,66]. The low-energy effective action for intersecting D-brane models has been given in [64,67,68] while explicit formulas for the soft-supersymmetry breaking terms used to generate the phenomenology in this work are contained in Ref. [12,69]. We assume that the gauge symmetry of the observable sector consist of only the MSSM below the usual GUT scale, $M_{\rm GUT} = 2.2 \cdot 10^{16}$ GeV, hence the observable supersymmetric phenomenology should remain consistent with formulae of Ref. [12]. To examine the dark matter content of the D6 model space, we investigate regions of the intersecting D6-brane model parameter space that satisfy all of the most current experimental constraints. The soft terms are input into MICROMEGAS 2.0.7 [70] using SUSPECT 2.34 [71] as a front end to run the soft terms down to the electroweak scale via the renormalization group equations (RGEs) and then to calculate the corresponding relic neutralino density, while μ is determined by the requirement of radiative electroweak symmetry breaking (REWSB). However, we do take $\mu > 0$ as suggested by the results of $g_\mu - 2$ for the muon. We use the current world average central value top quark mass of $m_t = 173.1$ GeV [72]. The direct-detection cross sections are calculated using MICROMEGAS 2.1 [73]. We apply the following experimental constraints:

(1) The 7 yr WMAP measurements of the cold dark matter density [74], $0.1088 \leq \Omega_\chi \leq 0.1158$. We also investigate another case where a neutralino LSP makes up a subdominant component and employ this possibility by removing the lower bound.
(2) The experimental limits on the flavor-changing neutral current (FCNC) process, $b \to s\gamma$. The results from the heavy flavor averaging group (HFAG) [75], in addition to the *BABAR*, Belle, and CLEO results, are: $Br(b \to s\gamma) = (355 \pm 24^{+9}_{-10} \pm 3) \times 10^{-6}$. There is also a more recent estimate [76] of $Br(b \to s\gamma) = (3.15 \pm 0.23) \times 10^{-4}$. For our analysis, we use the limits $2.86 \times 10^{-4} \leq Br(b \to s\gamma) \leq 4.18 \times 10^{-4}$, where experimental and theoretical errors are added in quadrature.
(3) The anomalous magnetic moment of the muon, $g_\mu - 2$. For this analysis we use the 2σ level boundaries, $11 \times 10^{-10} < a_\mu < 44 \times 10^{-10}$ [77].
(4) The process $B_s^0 \to \mu^+\mu^-$ where the decay has a $\tan^6\beta$ dependence. We take the upper bound to be $Br(B_s^0 \to \mu^+\mu^-) < 5.8 \times 10^{-8}$ [78].
(5) The LEP limit on the lightest *CP*-even Higgs boson mass, $m_h \geq 114$ GeV [79].

We present the updated WIMP-nucleon spin-independent cross-section contours in Figs. 1–5. Following the methodology in Ref. [12], we segregate the parameter space into distinctive scenarios of $m_{3/2}$ and $\tan\beta$. The five scenarios of $m_{3/2}$ and $\tan\beta$ first introduced in [12] and also analyzed here in this work were selected to be representative of a broad range of the experimentally allowed parameter space. To satisfy the LEP limit on the lightest *CP*-even Higgs boson mass, the gravitino mass needs to generally be a minimum of about $m_{3/2} = 500$ GeV, while the observability of the D6 model at the near-term LHC is questionable for a gravitino mass greater than $m_{3/2} = 700$ GeV due to very heavy superpartners, hence the range of $m_{3/2}$ used for these analyses. Likewise, the $\tan\beta$ examined here are those from near the minimum $\tan\beta$ possible to satisfy the experimental constraints, to a large $\tan\beta$ value representative of the high $\tan\beta$ region of the model space. The spin-independent contours in Figs. 1–5 represent the most current upper bounds from the CDMS [20] and XENON100 [21] experiments, as a function of the LSP mass. We find that the only regions significantly affected by the XENON100 constraints are those where the lightest neutralino and chargino are nearly degenerate. These points of chargino-neutralino degeneracy possess a nearly zero relic density, thus the neutralino would comprise only a tiny fraction of the total cold dark matter. Nonetheless, those regions with both chargino-neutralino and stau-neutralino coannihilation do subsist for potential supersymmetry and LSP discovery. Each figure is demarcated to clearly identify the appropriate areas of coannihilation.

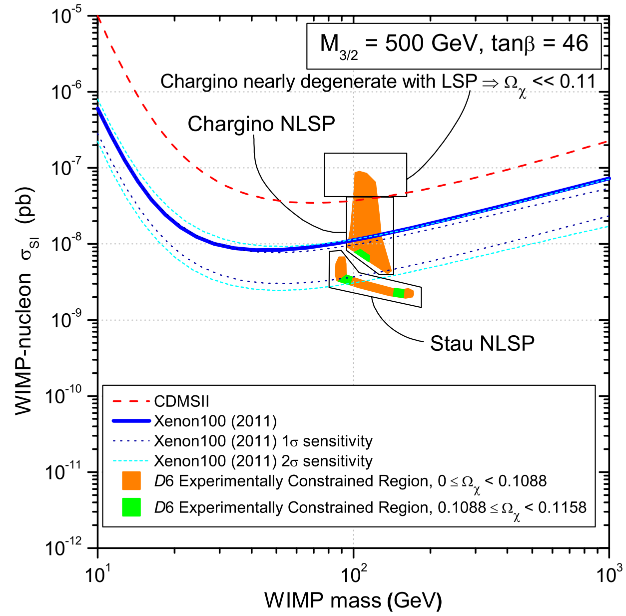

FIG. 1 (color online). Direct dark matter detection diagram associating the WIMP mass with the spin-independent annihilation cross-section $\sigma_{\rm SI}$. Delineated are the current upper bounds from the CDMS [20] and XENON100 [21] experiments. Shown is the experimentally viable parameter space for a gravitino mass $M_{3/2} = 500$ GeV and $\tan\beta = 46$. The boxes segregate the model space into the noted coannihilation regions.

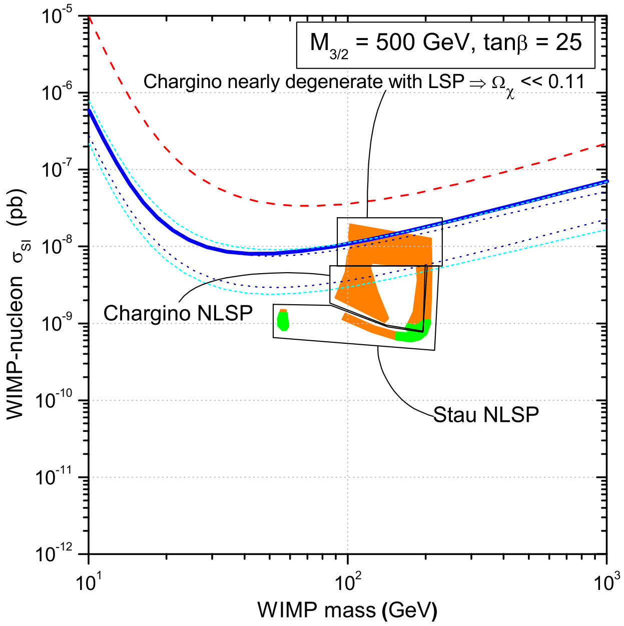

FIG. 2 (color online). Direct dark matter detection diagram associating the WIMP mass with the spin-independent annihilation cross-section σ_{SI}. Delineated are the current upper bounds from the CDMS [20] and XENON100 [21] experiments. Shown is the experimentally viable parameter space for a gravitino mass $M_{3/2} = 500$ GeV and $\tan\beta = 25$. The boxes segregate the model space into the noted coannihilation regions. See Fig. 1 for the legend describing the appropriate contours and regions.

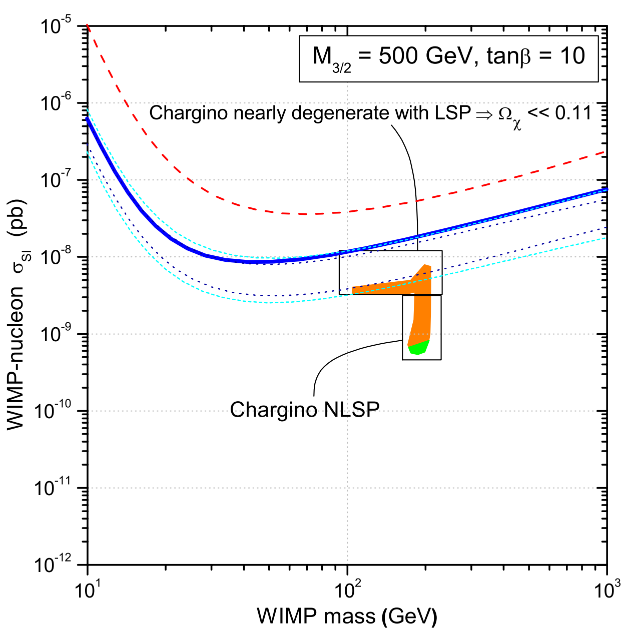

FIG. 3 (color online). Direct dark matter detection diagram associating the WIMP mass with the spin-independent annihilation cross-section σ_{SI}. Delineated are the current upper bounds from the CDMS [20] and XENON100 [21] experiments. Shown is the experimentally viable parameter space for a gravitino mass $M_{3/2} = 500$ GeV and $\tan\beta = 10$. The boxes segregate the model space into the noted coannihilation regions. See Fig. 1 for the legend describing the appropriate contours and regions.

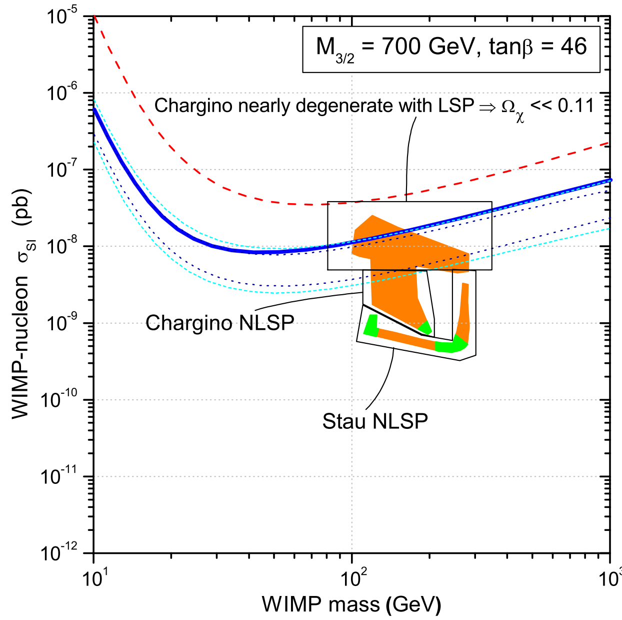

FIG. 4 (color online). Direct dark matter detection diagram associating the WIMP mass with the spin-independent annihilation cross-section σ_{SI}. Delineated are the current upper bounds from the CDMS [20] and XENON100 [21] experiments. Shown is the experimentally viable parameter space for a gravitino mass $M_{3/2} = 700$ GeV and $\tan\beta = 46$. The boxes segregate the model space into the noted coannihilation regions. See Fig. 1 for the legend describing the appropriate contours and regions.

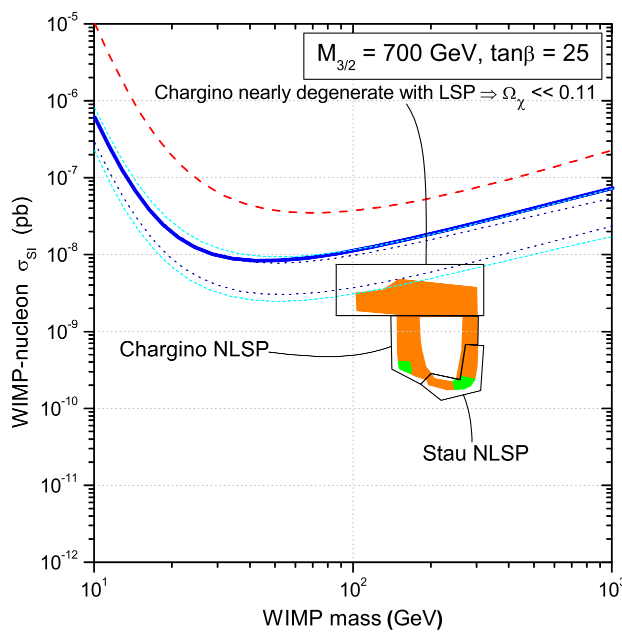

FIG. 5 (color online). Direct dark matter detection diagram associating the WIMP mass with the spin-independent annihilation cross-section σ_{SI}. Delineated are the current upper bounds from the CDMS [20] and XENON100 [21] experiments. Shown is the experimentally viable parameter space for a gravitino mass $M_{3/2} = 700$ GeV and $\tan\beta = 25$. The boxes segregate the model space into the noted coannihilation regions. See Fig. 1 for the legend describing the appropriate contours and regions.

V. CONCLUSION

We have discussed the presence of an extra nonanomalous $U(1)$ gauge symmetry in a realistic three-generation Pati-Salam model constructed with intersecting D6-branes in Type IIA string theory on a $T^6/(\mathbb{Z}_2 \times \mathbb{Z}_2)$ orientifold. As discussed in previous papers, the SM gauge couplings are unified at the string scale in this model, and it is possible to obtain realistic Yukawa matrices for quarks and leptons. Besides these favorable phenomenological features, we have shown that the additional $U(1)$ prohibits all dimension-4, 5, and 6 operators that mediate proton decay in the MSSM. In particular, we have shown that this $U(1)$ gives rise to a $U(1)_{3B+L}$ gauge symmetry once the Pati-Salam gauge symmetry is broken. Furthermore, it is possible to find linear combinations with $U(1)_{B-L}$ which lead to the effective promotion of baryon and lepton number to local gauge symmetries. Thus, the proton is essentially stable in this model.

As was mentioned, in the MSSM, rapid proton decay is allowed by the dimension-4 operators unless a discrete symmetry such as R-parity is imposed. However, even though these operators can be removed without invoking R-parity in the model considered here, R-parity can naturally surface from a spontaneously broken $U(1)_{B-L}$ provided certain conditions are satisfied which can be accomplished in the model. Thus, presuming the existence of R-parity giving rise to a stable LSP, we studied the direct-detection cross-sections for neutralino dark matter through application of the most current constraints from the XENON100 and CDMS experiments. We found that other than those regions with a lightest neutralino and chargino degeneracy, the D6 model space remains relatively intact and unaffected by the XENON100 constraints. These include regions of the parameter space where the relic density is generated through neutralino coannihilation with the stau and lightest chargino. However, this experiment will soon have sufficient reach to thoroughly test the model predictions for stable neutralino dark matter.

We should comment that it is really quite remarkable that a nonanomalous $U(1)$ gauge symmetry arises in the model which automatically leads to a stable proton with a very long lifetime. Given the observed long lifetime of the proton, proton stability can be considered one of the essential properties of any model of particle physics, string theory vacua, in particular, given the existence of the string landscape. It is known that models built on minimal $SU(5)$ do not satisfy this criteria as the dimension-5 operators which mediate proton decay are present in such constructions. However, in unified models based on $SU(5)$ such as flipped $SU(5)$, the proton may decay at a rate which is observable, but which satisfies current experimental limits. Although the model considered in this paper has a Pati-Salam structure, we have pointed out that it is not possible to obtain this Pati-Salam from a GUT such as $SO(10)$ due to the charges under the extra $U(1)$ carried by the matter supermultiplets. Thus, this scenario can be considered to give rise to a new 'GUT-less' paradigm where the proton is stable and the gauge couplings are unified at high energies, but where the gauge symmetry does not unify to a grand theory (other than of the Pati-Salam type).

ACKNOWLEDGMENTS

This was supported in part by the National Science Foundation under Grant No. PHY-0757394 (V. E. M), and by the Mitchell-Heep Chair in High Energy Physics and Department Of Energy Grant No. DE-FG03-95-Er-40917 (D. V. N. and J. A. M).

[1] P. Fayet, Nucl. Phys. **B90**, 104 (1975).
[2] P. Fayet, Phys. Lett. B **69**, 489 (1977).
[3] H.P. Nilles and S. Raby, Nucl. Phys. **B198**, 102 (1982).
[4] J.R. Ellis, J.S. Hagelin, D.V. Nanopoulos, and M. Srednicki, Phys. Lett. B **127**, 233 (1983); J.R. Ellis, J.S. Hagelin, D.V. Nanopoulos, K.A. Olive, and M. Srednicki, Nucl. Phys. **B238**, 453 (1984);
[5] H. Goldberg, Phys. Rev. Lett. **50**, 1419 (1983).
[6] R. Blumenhagen, M. Cvetic, P. Langacker, and G. Shiu, Annu. Rev. Nucl. Part. Sci. **55**, 71 (2005).
[7] R. Blumenhagen, B. Kors, D. Lust, and S. Stieberger, Phys. Rep. **445**, 1 (2007).
[8] M. Cvetic, T. Li, and T. Liu, Nucl. Phys. **B698**, 163 (2004).
[9] C.M. Chen, T. Li, and D.V. Nanopoulos, Nucl. Phys. **B740**, 79 (2006).
[10] C.M. Chen, T. Li, V.E. Mayes, and D.V. Nanopoulos, Phys. Lett. B **665**, 267 (2008).
[11] C.M. Chen, T. Li, V.E. Mayes, and D.V. Nanopoulos, Phys. Rev. D **77**, 125023 (2008).
[12] J.A. Maxin, V.E. Mayes, and D.V. Nanopoulos, Phys. Rev. D **81**, 015008 (2010).
[13] J.A. Maxin, V.E. Mayes, and D.V. Nanopoulos, Phys. Rev. D **79**, 123528 (2009).
[14] J.A. Maxin, V.E. Mayes, and D.V. Nanopoulos, arXiv:0912.5372.
[15] C.M. Chen, T. Li, V.E. Mayes, and D.V. Nanopoulos, J. Phys. G **35**, 095008 (2008).
[16] A.V. Belitsky, R.F. Lebed, and V.E. Mayes, Phys. Lett. B **697**, 343 (2011).
[17] R.F. Lebed and V.E. Mayes, arXiv:1103.4800.
[18] R.F. Lebed and V.E. Mayes,Phys. Rev. D **84**, 075016 (2011).

[19] J. R. Ellis, A. E. Faraggi, and D. V. Nanopoulos, Phys. Lett. B **419**, 123 (1998).
[20] Z. Ahmed *et al.* (The CDMS-II Collaboration), Science **327**, 1619 (2010).
[21] E. Aprile *et al.* (XENON100 Collaboration), Phys. Rev. D **84**, 052003 (2011).
[22] M. Cvetič, G. Shiu, and A. M. Uranga, Phys. Rev. Lett. **87**, 201801 (2001); Nucl. Phys. **B615**, 3 (2001).
[23] C. -M. Chen, G. V. Kraniotis, V. E. Mayes, D. V. Nanopoulos, and J. W. Walker, Phys. Lett. B **611**, 156 (2005).
[24] C. -M. Chen, G. V. Kraniotis, V. E. Mayes, D. V. Nanopoulos, and J. W. Walker, Phys. Lett. B **625**, 96 (2005).
[25] C. -M. Chen, V. E. Mayes, and D. V. Nanopoulos, Phys. Lett. B **633**, 618 (2006).
[26] L. E. Ibáñez and G. G. Ross, Nucl. Phys. **B368**, 3 (1992).
[27] A. Font, L. E. Ibáñez, and F. Quevedo, Phys. Lett. B **228**, 79 (1989).
[28] L. M. Krauss and F. Wilczek, Phys. Rev. Lett. **62**, 1221 (1989).
[29] S. P. Martin, Phys. Rev. D **46**, R2769 (1992).
[30] H. Murayama and A. Pierce, Phys. Rev. D **65**, 055009 (2002).
[31] S. M. Barr, Phys. Lett. B **112**, 219 (1982).
[32] J. P. Derendinger, J. E. Kim, and D. V. Nanopoulos, Phys. Lett. B **139**, 170 (1984).
[33] I. Antoniadis, J. R. Ellis, J. S. Hagelin, and D. V. Nanopoulos, Phys. Lett. B **194**, 231 (1987); **205**, 459 (1988); **208**, 209 (1988); **213**, 562 (1988); **231**, 65 (1989).
[34] R. Foot, G. C. Joshi, and H. Lew, Phys. Rev. D **40**, 2487 (1989).
[35] P. Fileviez Perez and M. B. Wise, Phys. Rev. D **82**, 011901 (2010).
[36] P. Fileviez Perez and M. B. Wise, Phys. Rev. D **84**, 055015 (2011).
[37] P. Fileviez Perez and M. B. Wise, J. High Energy Phys. 08 (2011) 068.
[38] M. Cvetic and P. Langacker, Phys. Rev. D **54**, 3570 (1996).
[39] R. Blumenhagen, M. Cvetič, S. Kachru, and T. Weigand, Annu. Rev. Nucl. Part. Sci. **59**, 269 (2009).
[40] J. L. Lopez and D. V. Nanopoulos, Phys. Rev. D **55**, 397 (1997).
[41] M. R. Buckley, D. Hooper, and J. L. Rosner, Phys. Lett. B **703**, 343 (2011).
[42] T. G. Rizzo, arXiv:hep-ph/0610104.
[43] P. Langacker, Rev. Mod. Phys. **81**, 1199 (2009).
[44] M. S. Carena, A. Daleo, B. A. Dobrescu, and T. M. P. Tait, Phys. Rev. D **70**, 093009 (2004).
[45] M. Bando, T. Kugo, K. Yamawaki, Phys. Rep. **164**, 217 (1988).
[46] (LEP, ALEPH, DELPHI, L3, OPAL, LEP Electroweak Working Group, SLD Electroweak Group, and SLD Heavy Flavor Group Collaborations), arXiv:hep-ex/0312023.
[47] A. Abulencia *et al.* (CDF Collaboration), Phys. Rev. Lett. **96**, 211802 (2006).
[48] T. Aaltonen *et al.* (CDF Collaboration), Phys. Rev. Lett. **102**, 091805 (2009).
[49] D. Acosta *et al.* (CDF Collaboration), Phys. Rev. Lett. **95**, 131801 (2005).
[50] K. Nakamura *et al.* (Particle Data Group Collaboration), J. Phys. G **37**, 075021 (2010).
[51] T. Aaltonen *et al.* (CDF Collaboration), Phys. Rev. D **79**, 112002 (2009).
[52] J. Alitti *et al.* (UA2 Collaboration), Nucl. Phys. **B400**, 3 (1993).
[53] P. J. Fox and E. Poppitz, Phys. Rev. D **79**, 083528 (2009).
[54] O. Adriani *et al.* (PAMELA Collaboration), Nature (London) **458**, 607 (2009).
[55] J. Chang *et al.* (ATIC Collaboration) Nature **456**, 362 (2008).
[56] M. R. Buckley, D. Hooper, J. Kopp, and E. Neil, Phys. Rev. D **83**, 115013 (2011).
[57] T. Aaltonen *et al.* (CDF Collaboration), Phys. Rev. D **83**, 112003 (2011).
[58] T. Aaltonen *et al.* (CDF Collaboration), Phys. Rev. Lett. **106**, 171801 (2011).
[59] S. Chatrchyan *et al.* (CMS Collaboration), J. High Energy Phys. 08 (2011) 005.
[60] V. M. Abazov *et al.* (D0 Collaboration), Phys. Rev. Lett. **107**, 011804 (2011).
[61] L. A. Anchordoqui, H. Goldberg, X. Huang, D. Lust, and T. R. Taylor, Phys. Lett. B **701**, 224 (2011).
[62] L. A. Anchordoqui, I. Antoniadis, H. Goldberg, X. Huang, D. Lust, and T. R. Taylor (unpublished).
[63] J. L. Hewett, T. G. Rizzo, arXiv:1106.0294.
[64] B. Kors and P. Nath, Nucl. Phys. **B681**, 77 (2004).
[65] J. A. Maxin, V. E. Mayes, and D. V. Nanopoulos, Phys. Rev. D **79**, 066010 (2009).
[66] J. A. Maxin, V. E. Mayes, and D. V. Nanopoulos, Phys. Lett. B **690**, 501 (2010).
[67] A. Brignole, L. E. Ibáñez, and C. Muñoz, in *Perspectives ion Supersymmetry* edited by G. L. Kane (World Scientific, Singapore, 1998) p. 125.
[68] D. Lüst, P. Mayr, R. Richter, and S. Stieberger, Nucl. Phys. **B696**, 205 (2004); A. Font and L. E. Ibanez, J. High Energy Phys. 03 (2005) 040.
[69] G. L. Kane, P. Kumar, J. D. Lykken, and T. T. Wang, Phys. Rev. D **71**, 115017 (2005);
[70] G. Belanger, F. Boudjema, A. Pukhov, and A. Semenov, Comput. Phys. Commun. **176**, 367 (2007); **174**, 577 (2006); **149**, 103 (2002).
[71] A. Djouadi, J. L. Kneur, and G. Moultaka, Comput. Phys. Commun. **176**, 426 (2007).
[72] (Tevatron Electroweak Working Group, CDF Collaboration, and D0 Collaboration), arXiv:0903.2503.
[73] G. Belanger, F. Boudjema, A. Pukhov, and A. Semenov, Comput. Phys. Commun. **180**, 747 (2009).
[74] D. N. Spergel *et al.* (WMAP Collaboration), Astrophys. J. Suppl. Ser. **148**, 175 (2003); **170**, 377 (2007); E. Komatsu *et al.*, Astrophys. J. Suppl. Ser. **192**, 18 (2011).

[75] E. Barberio *et al.* (Heavy Flavor Averaging Group (HFAG) Collaboration), arXiv:0704.3575.
[76] M. Misiak *et al.*, Phys. Rev. Lett. **98**, 022002 (2007).
[77] G. W. Bennett *et al.* (Muon g-2 Collaboration), Phys. Rev. Lett. **92**, 161802 (2004).
[78] T. Aaltonen *et al.* (CDF Collaboration), Phys. Rev. Lett. **100**, 101802 (2008).
[79] R. Barate *et al.* (LEP Working Group for Higgs boson searches and ALEPH Collaboration), Phys. Lett. B **565**, 61 (2003); W. M. Yao *et al.* (Particle Data Group), J. Phys. G **33**, 1 (2006).