Phenomenology of a three-family standardlike string model

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We discuss the phenomenology of a three-family supersymmetric standardlike model derived from the orientifold construction, in which the ordinary chiral states are localized at the intersection of branes at angles. In addition to the standard model group, there are two additional $U(1)'$ symmetries, one of which has family nonuniversal and therefore flavor changing couplings, and a quasihidden non-Abelian sector which becomes strongly coupled above the electroweak scale. The perturbative spectrum contains a fourth family of exotic [$SU(2)$-singlet] quarks and leptons, in which, however, the left-chiral states have unphysical electric charges. It is argued that these decouple from the low energy spectrum due to hidden sector charge confinement, and that anomaly matching requires the physical left-chiral states to be composites. The model has multiple Higgs doublets and additional exotic states. The moduli-dependent predictions for the gauge couplings are discussed. The strong coupling agrees with experiment for reasonable moduli, but the electroweak couplings are too small.

I. INTRODUCTION

Despite an enormous amount of promise and interest in superstring or M theory, it is still difficult to construct models with a fully realistic effective low energy four-dimensional limit from first principles. The (related) difficulties include the fact that only certain limiting corners of M theory are perturbative, there are many possible compactifications of the extra dimensions, the cosmological constant, the method of supersymmetry breaking, the difficulty of constructing a realistic spectrum for the effective theory, and the stabilization of moduli. One direction, the subject of this paper, is to examine concrete semirealistic constructions in as much detail as possible. One does not expect to find a fully realistic model. Rather, the goals are (1) to develop calculational techniques, (2) to suggest promising directions for new constructions, and (3) to find examples of possible new physics at the TeV scale that might emerge from explicit top-down string constructions. The latter are sometimes different from new physics motivated from bottom-up approaches. Of course, the features of a particular model may simply be shortcomings of that construction rather than generic predictions of classes of string theories. However, by examining promising constructions from a variety of corners of M theory one may obtain hints as to likely types of new TeV scale physics. Difficulties of this program are that most constructions are rather complicated, and that the unrealistic or unsuccessful aspects of each model often make it difficult to carry out detailed calculations.

For over a decade, there has been considerable effort in the construction of semirealistic string models in the framework of perturbative heterotic string theory [1]. In particular, a class of free-fermionic string models which contain the gauge group and matter content of the minimal supersymmetric standard model (MSSM) have been constructed [2,3].

These constructions have many interesting features, such as extended gauge structures and matter content [4]. Some of the features of one of these constructions are summarized in the Appendix.

The purpose of this paper is to examine the phenomenological issues in another calculable regime of M theory, namely, type II orientifolds. In recent years, the advent of D-branes has facilitated the construction of semirealistic string models using conformal field theory techniques, as illustrated by the various four-dimensional $\mathcal{N}=1$ supersymmetric type II orientifolds ([5–16] and references therein). A promising direction to obtain chiral theories is by constructing models with D-branes intersecting at angles [17]. This fact (or its T-dual version, i.e., branes with flux) has been exploited in [18–20] to construct semirealistic string models (see also [21]). However, the constraints on supersymmetric four-dimensional models are rather restrictive. Despite the remarkable progress in developing techniques of orientifold constructions, there is only one orientifold model [15,16] that has been constructed so far with the ingredients of the MSSM: $\mathcal{N}=1$ supersymmetry, the standard model (SM) gauge group as a part of the gauge structure, and candidate fields for the three generations of quarks and leptons as well as the electroweak Higgs doublets. We hope that by studying the phenomenology of this model in detail, we can probe some of the generic features and predictions of string models derived from the orientifold approach.

In this paper we concentrate on direct compactifications of the underlying M theory to a four-dimensional field theory containing the MSSM (i.e., without having an intermediate four-dimensional grand unified theory). We focus on the case in which the fundamental scale is comparable to the Planck scale, i.e., the case with no very large extra dimensions.

In Sec. II we briefly summarize the construction of the model, including the gauge factors and the quantum numbers

of the chiral and nonchiral spectrum. The properties of the perturbative spectrum are discussed in more detail in Sec. III, including the properties of the multiple Higgs doublets, the three regular families, the fourth exotic family, and alternative assignments. The properties of the additional $U(1)'$ gauge interactions and the possibilities for breaking them at the electroweak or intermediate scales are discussed in Sec. IV. Section V is concerned with the gauge couplings. The model does not have the conventional form of gauge unification because each group factor is associated with a different set of branes. However, the string-scale couplings are predicted in terms of the ratio of the Planck to string scales and a geometric factor. The low energy electroweak couplings are too small due to the multiple Higgs fields and exotic matter, while the strong coupling is more reasonable. The quasihidden sector groups are asymptotically free. The implications of these results for the spectrum are described in Sec. VI. In particular, the fractionally charged exotic states presumably disappear from the low energy spectrum due to hidden sector charge confinement, to be replaced by composite states with the appropriate quantum numbers to form the left-handed components of an exotic fourth family. The results are summarized and contrasted with a particular heterotic construction in Sec. VII. A more detailed description of the features of that construction is given in the Appendix. A detailed discussion of the Yukawa couplings of the model will be presented separately [22], and further implications of the strong couplings for moduli stabilization and supersymmetry breaking is under investigation [23].

II. DESCRIPTION OF THE MODEL

For completeness, let us describe the construction of the model in [15], which is obtained by compactifying type IIA string theory on a $\mathbf{T}^6/(\mathbf{Z}_2 \times \mathbf{Z}_2)$ orientifold. We considered a general framework in which the D-branes are not necessarily parallel to the orientifold planes, which gives rise to new possibilities of embedding the gauge sector in the same background geometry. (In the T-dual picture [6], this corresponds to turning on a background flux of the gauge fields on the D-branes.)

The generators θ, ω of $\mathbf{Z}_2 \times \mathbf{Z}_2$ act as θ: $(z_1, z_2, z_3) \rightarrow (-z_1, -z_2, z_3)$, and ω: $(z_1, z_2, z_3) \rightarrow (z_1, -z_2, -z_3)$ on the complex coordinates z_i of $\mathbf{T}^6 = \mathbf{T}^2 \times \mathbf{T}^2 \times \mathbf{T}^2$. The orientifold action is ΩR, where Ω is world-sheet parity, and R acts by R: $(z_1, z_2, z_3) \rightarrow (\bar{z}_1, \bar{z}_2, \bar{z}_3)$. The model contains four kinds of O6-planes, associated with the actions of ΩR, $\Omega R \theta$, $\Omega R \omega$, and $\Omega R \theta \omega$. The closed string sector contains gravitational supermultiplets as well as orbifold moduli and is straightforward to determine, and so in what follows we will focus on the open string (charges) spectrum. The cancellation of the Ramond-Ramond (RR) tadpoles from the orientifold planes requires an introduction of K stacks of N_a $D6$-branes ($a=1,...,K$) wrapped on three-cycles $[\Pi_a]$ (which for simplicity is taken to be the product of 1-cycles of the two-tori):

$$[\Pi_a] = \prod_{i=1}^{3} (n_a^i [a_i] + m_a^i [b_i]) \qquad (1)$$

and similarly for their orientifold images under ΩR. The cycles that the $D6$-branes and their orientifold images wrap around are specified by the wrapping numbers (n_a^i, m_a^i). The number of $D6$-branes and the associated wrapping numbers are constrained by (i) cancellation of RR tadpoles and (ii) supersymmetry.

Analogous to the situation in [20], models with all tori orthogonal lead to an even number of families. Hence, as in [20], we consider models with one tilted T^2, where the tilting parameter is discrete and has a unique nontrivial value [24]. This mildly modifies the closed string sector, but has an important impact on the open string sector. Namely, a D-brane 1-cycle (n_a^i, m_a^i) along a tilted torus is mapped to $(n^i, -m^i - n^i)$. It is convenient to define $\tilde{m}^i = m^i + \frac{1}{2} n^i$, and label the cycles as $(n^i, \tilde{m}^i)$.

Let us summarize the results for $D6$-branes not parallel to O6-planes (for zero angles, the spectrum follows from [6]). The aa sector (strings stretched within a single stack of $D6_a$-branes) is invariant under θ, ω, and is exchanged with $a'a'$ by the action of ΩR. For the gauge group, the θ projection breaks $U(N_a)$ to $U(N_a/2) \times U(N_a/2)$, and ω identifies both factors, leaving $U(N_a/2)$. Concerning the matter multiplets, we obtain three adjoint $N=1$ chiral multiplets.

The $ab+ba$ sector (strings stretched between $D6_a$- and $D6_b$-branes) is invariant, as a whole, under the orbifold projections, and is mapped to the $b'a' + a'b'$ sector by ΩR. The matter content before any projection would be given by I_{ab} chiral fermions in the bifundamental $(\square_a, \bar{\square}_b)$ of $U(N_a) \times U(N_b)$, where

$$I_{ab} = (n_a^1 m_b^1 - n_b^1 m_a^1)(n_a^2 m_b^2 - n_b^2 m_a^2)$$
$$\times (n_a^3 \tilde{m}_b^3 - n_b^3 \tilde{m}_a^3)$$
$$\equiv \prod_{i=1}^{3} I_{ab}^i$$

is the intersection number of the wrapped cycles (see [18,19]), and the sign of I_{ab} denotes the chirality of the corresponding fermion ($I_{ab} < 0$ gives left-handed fermions in our convention). For supersymmetric intersections, additional massless scalars complete the corresponding chiral supermultiplet. In principle, one needs to take into account the orbifold action on the intersection point. However the final result turns out to be insensitive to this subtlety and is still given by I_{ab} chiral multiplets in the $(\square_a, \bar{\square}_b)$ of $U(N_a/2) \times U(N_b/2)$. A similar effect takes place in the $ab' + b'a$ sector, for $a \neq b$, where the final matter content is given by $I_{ab'}$ chiral multiplets in the bifundamental $(\square_a, \square_b)$.

For the $aa' + a'a$ sector the orbifold action on the intersection points turns out to be crucial. For intersection points invariant under the orbifold, the orientifold projection leads to a two-index antisymmetric representation of $U(N_a/2)$, except for states with θ and ω eigenvalue $+1$, where it yields a two-index symmetric representation. For points not fixed under some orbifold element, say two points fixed under ω and exchanged by θ, one simply keeps one point, and does not

TABLE I. $D6$-brane configuration for the three-family model.

Type	N_a	$(n_a^1, m_a^1) \times (n_a^2, m_a^2) \times (n_a^3, \tilde{m}_a^3)$	Group
A_1	8	$(0,1) \times (0,-1) \times (2,\tilde{0})$	$Q_{8,8'}$
A_2	2	$(1,0) \times (1,0) \times (2,\tilde{0})$	$Sp(2)_A$
B_1	4	$(1,0) \times (1,-1) \times (2,\widetilde{3/2})$	$SU(2)$
B_2	2	$(1,0) \times (0,1) \times (0,\widetilde{-1})$	$Sp(2)_B$
C_1	6+2	$(1,-1) \times (1,0) \times (1,\widetilde{1/2})$ i	$SU(3), Q_3, Q_1$
C_2	4	$(0,1) \times (1,0) \times (0,\widetilde{-1})$	$Sp(4)$

impose the ω projection. Equivalently, one considers all possible eigenvalues for ω, and applies the above rule to read off whether the symmetric or the antisymmetric survives. A closed formula for the chiral piece in this sector can be found in [15].

In addition to the chiral multiplets, there can be vector-like multiplets from the $ab+ba$, $ab'+b'a$, and $aa'+a'a$ sectors. This happens when $I_{ab}^i=0$ for a single $i=i_0$, the intersection number is zero. Instead of having only left-handed or right-handed chiral multiplets at the intersection, both the left-handed and the right-handed chiral multiplets are present. The multiplicity of vector-like multiplets is given by

$$I_{ab} = \prod_{i \neq i_0} I_{ab}^i . \tag{2}$$

These vector-like multiplets are generically massive as explained in [15].

The condition that the system of branes preserve $N=1$ supersymmetry requires [17] that each stack of $D6$-branes is related to the O6-planes by a rotation in $SU(3)$: denoting by θ_i the angles the $D6$-brane forms with the horizontal direction in the ith two-torus, supersymmetry preserving configurations must satisfy $\theta_1 + \theta_2 + \theta_3 = 0$. In order to simplify the supersymmetry conditions within our search for realistic models, we considered a particular ansatz: $(\theta_1, \theta_2, 0)$, $(\theta_1, 0, \theta_3)$, or $(0, \theta_2, \theta_3)$.

Due to the smaller number of O6-planes in tilted configurations, the RR tadpole conditions are very stringent for more than one tilted torus. Focusing on tilting just the third torus, the search for theories with $U(3)$ and $U(2)$ gauge factors carried by branes at angles and three left-handed quarks, turns out to be very constraining, at least within our ansatz. We have found essentially a unique solution. The $D6$-brane configuration with wrapping numbers $(n_a^i, \tilde{m}_a^i)$ is given in Table I.

The eight $D6$-branes labeled C_1 are split in two parallel but not overlapping stacks of 6 and 2 branes, and hence lead to a gauge group $U(3) \times U(1)$. Interestingly, a linear combination of the two $U(1)$'s is actually a generator within the $SU(4)$ arising for coincident branes. This ensures that this $U(1)$ is automatically nonanomalous and massless (free of

linear couplings to untwisted moduli) [19,25], and turns out to be crucial in the appearance of hypercharge in this model.

For convenience we consider the 8 $D6$-branes labeled A_1 to be away from the O6-planes in all three complex planes. This leads to two $D6$-branes that can move independently [hence give rise to a group $U(1)^2$], plus their θ, ω, and ΩR images. These $U(1)$'s are also automatically nonanomalous and massless. In the effective theory, this corresponds to Higgsing of $USp(8)$ down to $U(1)^2$.

The surviving non-Abelian gauge group is $SU(3) \times SU(2) \times Sp(2) \times Sp(2) \times Sp(4)$. The $SU(3) \times SU(2)$ corresponds to the MSSM, while the last three factors form a quasihidden sector (i.e., most states are charged under one sector or the other, but there are a few which couple to both). In addition, there are three nonanomalous $U(1)$ factors and two anomalous ones. The generators Q_3, Q_1, and Q_2 refer to the $U(1)$ factor within the corresponding $U(n)$, while Q_8, Q_8' are the $U(1)$'s arising from the $USp(8)$. $Q_3/3$ and Q_1 are essentially the baryon (B) and lepton (L) number, respectively, while $(Q_8+Q_8')/2$ is analogous to the generator T_{3R} occurring in left-right symmetric extensions of the SM. The hypercharge is defined as

$$Q_Y = \frac{1}{6}Q_3 - \frac{1}{2}Q_1 + \frac{1}{2}(Q_8 + Q_8'). \tag{3}$$

From the above comments, Q_Y as defined guarantees that $U(1)_Y$ is massless. There are two additional surviving nonanomalous $U(1)$'s, i.e., $B-L = Q_3/3 - Q_1$ and $Q_8 - Q_8'$. The gauge bosons corresponding to the anomalous $U(1)$ generators $B+L$ and Q_2 acquire string-scale masses, so those generators act like global symmetries on the effective four-dimensional theory.

The spectrum of chiral multiplets in the open string sector is tabulated in Table II. For completeness, we also give the spectrum of the vector-like multiplets in Tables III–V even though these multiplets are generically massive. The theory contains three standard model families, multiple Higgs candidates, a number of exotic chiral (but anomaly-free) fields, and multiplets which transform as adjoints or singlets under the SM gauge group. The spectrum is discussed in more detail in the following section.

III. THE PERTURBATIVE SPECTRUM

In this section we describe the simplest version of the perturbative chiral spectrum. The modifications expected due to strong coupling effects in the quasihidden sector will be described in Sec. VI.

The A_1B_1 sector contains 24 Higgs doublets, i.e., 12 H_U, H_D pairs, where H_U (H_D) has the appropriate hypercharge to generate masses for the charge 2/3 quarks and neutrinos (the charge $-1/3$ quarks and charged leptons). Six pairs have Q_8 charges, and the other six $Q_{8'}$ charges. The existence of so many doublets is the most unrealistic feature of the model,

TABLE II. The chiral spectrum of the open string sector in the three-family model. To be complete, we also list (in the bottom part of the table, below the horizontal line) the chiral states from the aa sectors, which are not localized at the intersections.

Sector	$SU(3)\times SU(2)\times Sp(2)_B\times Sp(2)_A\times Sp(4)$	Q_3	Q_1	Q_2	Q_8	Q_8'	Q_Y	Q_8-Q_8'	Field
A_1B_1	$3\times2\times(1,\bar{2},1,1,1)$	0	0	−1	±1	0	$\pm\frac{1}{2}$	±1	H_U, H_D
	$3\times2\times(1,\bar{2},1,1,1)$	0	0	−1	0	±1	$\pm\frac{1}{2}$	∓1	H_U, H_D
A_1C_1	$2\times(\bar{3},1,1,1,1)$	−1	0	0	±1	0	$\frac{1}{3},-\frac{2}{3}$	1, −1	$\bar{D}, \bar{U}$
	$2\times(\bar{3},1,1,1,1)$	−1	0	0	0	±1	$\frac{1}{3},-\frac{2}{3}$	−1, 1	$\bar{D}, \bar{U}$
	$2\times(1,1,1,1,1)$	0	−1	0	±1	0	1, 0	1, −1	$\bar{E}, \bar{N}$
	$2\times(1,1,1,1,1)$	0	−1	0	0	±1	1, 0	−1, 1	$\bar{E}, \bar{N}$
B_1C_1	$(3,\bar{2},1,1,1)$	1	0	−1	0	0	$\frac{1}{6}$	0	Q_L
	$(1,\bar{2},1,1,1)$	0	1	−1	0	0	$-\frac{1}{2}$	0	L
B_1C_2	$(1,2,1,1,4)$	0	0	1	0	0	0	0	
B_2C_1	$(3,1,2,1,1)$	1	0	0	0	0	$\frac{1}{6}$	0	
	$(1,1,2,1,1)$	0	1	0	0	0	$-\frac{1}{2}$	0	
B_1C_1'	$2\times(3,2,1,1,1)$	1	0	1	0	0	$\frac{1}{6}$	0	Q_L
	$2\times(1,2,1,1,1)$	0	1	1	0	0	$-\frac{1}{2}$	0	L
B_1B_1'	$2\times(1,1,1,1,1)$	0	0	−2	0	0	0	0	
	$2\times(1,3,1,1,1)$	0	0	2	0	0	0	0	
A_1A_1	$3\times8\times(1,1,1,1,1)$	0	0	0	0	0	0	0	
	$3\times4\times(1,1,1,1,1)$	0	0	0	±1	±1	±1	0	
	$3\times4\times(1,1,1,1,1)$	0	0	0	±1	∓1	0	±2	
	$3\times(1,1,1,1,1)$	0	0	0	±2	0	±1	±2	
	$3\times(1,1,1,1,1)$	0	0	0	0	±2	±1	∓2	
A_2A_2	$3\times(1,1,1,1,1)$	0	0	0	0	0	0	0	
B_1B_1	$3\times(1,3,1,1,1)$	0	0	0	0	0	0	0	
	$3\times(1,1,1,1,1)$	0	0	0	0	0	0	0	
B_2B_2	$3\times(1,1,1,1,1)$	0	0	0	0	0	0	0	
C_1C_1	$3\times(8,1,1,1,1)$	0	0	0	0	0	0	0	
	$3\times(1,1,1,1,1)$	0	0	0	0	0	0	0	
C_2C_2	$3\times(1,1,1,1,5+1)$	0	0	0	0	0	0	0	

and is a major cause for the unrealistic predictions for the low energy gauge couplings in the MSSM sector. Nevertheless, it at least suggests the possibility of additional Higgs doublets, and the associated phenomenological consequences, such as a richer Higgs-neutralino-chargino spectrum and Higgs-mediated flavor changing or CP violation effects. The problem for this construction is aggravated by the fact that there is no satisfactory mechanism to generate effective supersymmetric masses (μ terms) for all of the Higgs multiplets. Elementary μ terms are not generated in the string construction (and are also forbidden by Q_2). An effective μ would be generated by the vacuum expectation value (VEV) of the scalar component of a superfield S if there were a term SH_UH_D in the superpotential [26,27]. However, there are no $SU(2)$-singlet, $Q_Y=0$ states in the chiral spectrum with the gauge quantum numbers to play the role of S. In particular, the $\bar{N}H_UH_D$ couplings are forbidden by Q_1 and Q_2. However, there are nonchiral B_1B_1' states with the appropriate quantum numbers. There are couplings between the A_1B_1, B_1B_1', and $B_1'A_1$ sectors. Since the $B_1'A_1$

sector is the same as $A_1'B_1$, and furthermore, $A_1=A_1'$ (because the A_1 brane is the same as its own orientifold image), in principle, there are nonzero couplings of the form $(A_1B_1)(B_1B_1')(A_1B_1)$. Hence if these nonchiral states are not massive, they can play the role of the S field. An explicit calculation of the above coupling from the corresponding string amplitude, however, is necessary to verify that there are no additional stringy symmetries to forbid such couplings.

The B_1C_1' intersection yields two families of quark and lepton doublets, while a third is associated with B_1C_1. Only the first two can have Yukawa couplings to the $H_{U,D}$ because of Q_2, so one fermion family will remain massless. (The Yukawa couplings of the model are discussed in detail in [22].) The A_1C_1 sector contains four families of $SU(2)$-singlet antiquarks and antileptons, $\bar{U}$, $\bar{D}$, $\bar{E}$, and $\bar{N}$, two charged under Q_8 and two under $Q_{8'}$. Three of these should be the partners of the quark and lepton doublets.

The B_1C_2 and B_2C_1 states couple to both the ordinary and quasihidden sector gauge groups, and they carry the frac-

TABLE III. Nonchiral spectrum from the *ab* sectors in the three-family model. These states, as well as those in Tables IV and V, are generically massive.

Sector	$SU(3) \times SU(2) \times Sp(2)_B \times Sp(2)_A \times Sp(4)$	Q_3	Q_1	Q_2	Q_8	Q'_8	Q_Y	$Q_8 - Q'_8$
A_1A_2	$2 \times (1,1,1,2,1) + 2 \times (1,1,1,\bar{2},1)$	0	0	0	± 1	0	$\pm \frac{1}{2}$	± 1
	$2 \times (1,1,1,2,1) + 2 \times (1,1,1,\bar{2},1)$	0	0	0	0	± 1	$\pm \frac{1}{2}$	∓ 1
A_1B_2	$2 \times 2 \times (1,1,2,1,1) + 2 \times 2 \times (1,1,\bar{2},1,1)$	0	0	0	± 1	0	$\pm \frac{1}{2}$	± 1
	$2 \times 2 \times (1,1,2,1,1) + 2 \times 2 \times (1,1,\bar{2},1,1)$	0	0	0	0	± 1	$\pm \frac{1}{2}$	∓ 1
A_1C_2	$2 \times 2 \times (1,1,1,1,4) + 2 \times 2 \times (1,1,1,1,\bar{4})$	0	0	0	± 1	0	$\pm \frac{1}{2}$	± 1
	$2 \times 2 \times (1,1,1,1,4) + 2 \times 2 \times (1,1,1,1,\bar{4})$	0	0	0	0	± 1	$\pm \frac{1}{2}$	∓ 1
A_2B_1	$3 \times (1,2,1,\bar{2},1)$	0	0	1	0	0	0	0
	$3 \times (1,\bar{2},1,2,1)$	0	0	-1	0	0	0	0
A_2B_2	$2 \times (1,1,2,\bar{2},1) + 2 \times (1,1,\bar{2},2,1)$	0	0	0	0	0	0	0
A_2C_1	$(3,1,1,\bar{2},1)$	1	0	0	0	0	$\frac{1}{6}$	0
	$(\bar{3},1,1,2,1)$	-1	0	0	0	0	$-\frac{1}{6}$	0
	$(1,1,1,\bar{2},1)$	0	1	0	0	0	$-\frac{1}{2}$	0
	$(1,1,1,2,1)$	0	-1	0	0	0	$\frac{1}{2}$	0
A_2C_2	$2 \times (1,1,1,2,\bar{4}) + 2 \times (1,1,1,\bar{2},4)$	0	0	0	0	0	0	0
B_1B_2	$(1,2,\bar{2},1,1)$	0	0	1	0	0	$-\frac{1}{2}$	0
	$(1,\bar{2},2,1,1)$	0	0	-1	0	0	$\frac{1}{2}$	0
B_2C_2	$(1,1,2,1,\bar{4}) + (1,1,\bar{2},1,4)$	0	0	0	0	0	0	0
C_1C_2	$(3,1,1,1,\bar{4})$	1	0	0	0	0	$\frac{1}{6}$	0
	$(\bar{3},1,1,1,4)$	-1	0	0	0	0	$-\frac{1}{6}$	0
	$(1,1,1,1,\bar{4})$	0	1	0	0	0	$-\frac{1}{2}$	0
	$(1,1,1,1,4)$	0	-1	0	0	0	$\frac{1}{2}$	0

tional electric charges $\pm \frac{1}{2}, \frac{1}{6}$. In particular, the B_2C_1 states include two $SU(3)$ triplets and two $SU(3)$ singlets. All are $SU(2)$ singlets. These would be the natural partners of the extra $\bar{U}$, $\bar{D}$, $\bar{E}$, and $\bar{N}$ family from A_1C_1 except that they have the wrong electric and hypercharges.[1] It will be argued in Sec. VI that these states most likely disappear from the physical low energy spectrum due to charge confinement for the strongly coupled hidden sector gauge groups, to be replaced with composite fields with the appropriate quantum numbers to be the [$SU(2)$-singlet] partners of the fourth family of $\bar{U}$, $\bar{D}$, $\bar{E}$, and $\bar{N}$.

The $B_1B'_1$ and *aa* sector chiral supermultiplets in Table II include three $SU(3)$ octets and five $SU(2)$ triplets, as well as a number of non-Abelian singlets with nonzero hyper-

charge. The *aa* states are not localized at intersections, and there is no known mechanism to give them large masses. Hence they contribute significantly to the running of the gauge couplings. (We will actually ignore them, hoping that some mass mechanism will be eventually found.) These kinds of states (adjoints of the non-Abelian gauge symmetry) are a generic problem for orientifold constructions with intersecting branes; since each brane wraps a (flat) supersymmetric three-cycle of the six-torus, these three-cycles are *not* rigid and thus the adjoint matter corresponds to the moduli associated with the translational invariance of each three-cycle. One possibility to remove these states would be to put branes on curved (Calabi-Yau) space with rigid three-cycles, which is beyond the scope of the original construction of the model. The sector also contains a number of SM singlets, and three $Sp(4)$ 5-plets.

One can consider alternative identifications of some of the multiplets, due to the fact that the lepton doublets L and the Higgs doublets H_D have the same SM quantum numbers. For example, one could interpret one or both of the $SU(3)$ singlets in $B_1C'_1$ to be H'_D states, and some of the doublets in A_1B_1 as leptons. One could give masses to all of the charged leptons in this way, but not all of the quarks. Also, if the scalar component of $\bar{N}$ acquired a VEV, some, but not all, of the Higgs fields could have effective μ terms. We do not pursue this alternative further.

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[1]We have explored the possibility of using an alternative hypercharge definition, in which Q_Y includes as an additional term the diagonal generator of the first $Sp(2)$ group. However, the required $Sp(2)$ breaking cannot occur without breaking supersymmetry because in this sector there are only two branes and thus there are not enough branes available to split the branes in a $\mathbf{Z}_2 \times \mathbf{Z}_2$ invariant manner and be compatible with the orientifold projection. Breaking of this $Sp(2)$ symmetry by some (perhaps nonperturbative) mechanism at the TeV scale might still be a possibility, but we concentrate here on the alternative mechanism described in Sec. VI.

TABLE IV. Nonchiral spectrum from the ab' sectors in the three-family model.

Sector	$SU(3)\times SU(2)\times Sp(2)_B\times Sp(2)_A\times Sp(4)$	Q_3	Q_1	Q_2	Q_8	Q'_8	Q_Y	$Q_8-Q'_8$
$A_1B'_2$	$2\times2\times(1,1,2,1,1)+2\times2\times(1,1,\bar{2},1,1)$	0	0	0	±1	0	$\pm\frac{1}{2}$	±1
	$2\times2\times(1,1,2,1,1)+2\times2\times(1,1,\bar{2},1,1)$	0	0	0	0	±1	$\pm\frac{1}{2}$	∓1
$A_1C'_2$	$2\times2\times(1,1,1,1,4)+2\times2\times(1,1,1,1,\bar{4})$	0	0	0	±1	0	$\pm\frac{1}{2}$	±1
	$2\times2\times(1,1,1,1,4)+2\times2\times(1,1,1,1,\bar{4})$	0	0	0	0	±1	$\pm\frac{1}{2}$	∓1
$A_2B'_1$	$3\times(1,2,1,2,1)$	0	0	1	0	0	0	0
	$3\times(1,\bar{2},1,\bar{2},1)$	0	0	-1	0	0	0	0
$A_2B'_2$	$2\times(1,1,2,2,1)+2\times(1,1,\bar{2},\bar{2},1)$	0	0	0	0	0	0	0
$A_2C'_1$	$(3,1,1,2,1)$	1	0	0	0	0	$\frac{1}{6}$	0
	$(\bar{3},1,1,\bar{2},1)$	-1	0	0	0	0	$-\frac{1}{6}$	0
	$(1,1,1,2,1)$	0	1	0	0	0	$-\frac{1}{2}$	0
	$(1,1,1,\bar{2},1)$	0	-1	0	0	0	$\frac{1}{2}$	0
$A_2C'_2$	$2\times(1,1,1,2,4)+2\times(1,1,1,\bar{2},\bar{4})$	0	0	0	0	0	0	0
$B_1B'_2$	$(1,2,2,1,1)$	0	0	1	0	0	0	0
	$(1,\bar{2},\bar{2},1,1)$	0	0	-1	0	0	0	0
$B_2C'_2$	$(1,1,2,1,3)+(1,1,\bar{2},1,\bar{4})$	0	0	0	0	0	0	0
$C_1C'_2$	$(3,1,1,1,4)$	1	0	0	0	0	$\frac{1}{6}$	0
	$(\bar{3},1,1,1,\bar{4})$	-1	0	0	0	0	$-\frac{1}{6}$	0
	$(1,1,1,1,4)$	0	1	0	0	0	$-\frac{1}{2}$	0
	$(1,1,1,1,\bar{4})$	0	-1	0	0	0	$\frac{1}{2}$	0

IV. THE ADDITIONAL $U(1)$ FACTORS

As described in Sec. II, there are two surviving nonanomalous $U(1)$'s, $B-L=Q_3/3-Q_1$ and $Q_8-Q'_8$, in addition to hypercharge.[2] The $Q_8-Q'_8$ charges are not family universal for the antiquarks and antileptons. For example, two of the $\bar{U}$ have $Q_8-Q'_8=-1$ and two have $+1$. When family mixing is considered, there may be flavor changing neutral current (FCNC) couplings of the corresponding Z boson to the $\bar{U}$ quarks [28], provided that the two types of $\bar{U}$'s mix with each other. (This will occur provided there are H_U scalars of both types, i.e., with $Q_8\neq0$ and with $Q_{8'}\neq0$, with nonzero VEVs.) Similar statements apply to the $\bar{D}$, $\bar{E}$, and $\bar{N}$ couplings. Such couplings could lead to decays such as $B_s\rightarrow\mu^+\mu^-$ or $\tau\rightarrow\mu e^+e^-$. On the other hand, the $B-L$ charges are family universal (even though one of the three families of Q and L has a different origin than the other two), so there are no flavor-violating $B-L$ couplings.

Limits on the mass and mixings of extra Z' bosons depend on their couplings, but typically $M_{Z'}>500-800$ GeV, and the Z-Z' mixing angle is less than few $\times10^{-3}$ [29]. One possibility is for a $U(1)'$ to be broken at a large scale intermediate between the TeV and Planck scales. This can occur if there are two or more scalar fields with opposite signs for their $U(1)'$ charges, so that the symmetry breaking is along a D (and F)-flat direction [30]. This would be desirable for implementing a neutrino seesaw mechanism if the heavy singlet neutrino carries a nonzero $U(1)'$ charge. Alternatively, the breaking can occur at the TeV scale or lower [27], in which case the minimum need not be supersymmetry conserving. In either case, one expects to first generate masses for and possibly mixing between the two new Z' bosons via the VEVs of SM-singlet fields that are large compared to the electroweak scale. A small mixing of these states with the Z would then be induced by Higgs doublet VEVs if they carry $U(1)'$ charges.[3]

The only SM-singlet fields which couple to the extra $U(1)$''s are: (1) the two pairs of $\bar{N}$ states with $B-L=1$ and $Q_8=-1$ or $Q'_8=-1$; we refer to these as $\bar{N}(8)$ and $\bar{N}(8')$, respectively. (2) The $Q_Y=B-L=0$ and $Q_8-Q'_8=\pm2$ states in the (otherwise problematic) A_1A_1 sector. The A_1A_1 states, if present in the spectrum, could lead to a D-flat direction for which $Q_8-Q'_8$ (but not $B-L$) is broken at an intermediate

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[2]Strictly speaking, the three independent $U(1)'$ charges are $Q_8-Q'_8$, $Q_8+Q'_8$, and $B-L$. One can rotate the corresponding gauge bosons so that they couple to $Q_8-Q'_8$, Q_Y, and $\beta(Q_8+Q'_8)-1/\beta(B-L)$, where $\beta\equiv g_{8+}/g_{B-L}$ is the ratio of $Q_8+Q'_8$ and $B-L$ gauge couplings.

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[3]In a multi-Higgs model, one can in principle generate the entire Z-Z' mass matrix by the VEVs of Higgs doublets. (One of the mass eigenvalues vanishes for a single doublet.) The mixing can vanish for one or more $H_{U,D}$ pairs (along a D-flat direction for more than one pair). However, the large Z'/Z mass ratio would require a very large Z' coupling to the doublets, which is not the case here.

TABLE V. Nonchiral spectrum from the aa' sectors in the three-family model.

Sector	$SU(3)\times SU(2)\times Sp(2)_B\times Sp(2)_A\times Sp(4)$	Q_3	Q_1	Q_2	Q_8	Q_8'	Q_Y	Q_8-Q_8'
B_1B_1'	$4\times(1,1,1,1,1)$	0	0	±2	0	0	0	0
C_1C_1'	$2\times(3,1,1,1,1)$	2	0	0	0	0	$\frac{1}{3}$	0
	$2\times(\bar{3},1,1,1,1)$	-2	0	0	0	0	$-\frac{1}{3}$	0

scale. Depending on whether these A_1A_1 states are relevant, one could have either that Q_8-Q_8' is broken at an intermediate scale and $B-L$ at the TeV scale [by $\bar{N}(8)$ and $\bar{N}(8')$]; or that both are broken at the TeV scale by $\bar{N}(8)$ and $\bar{N}(8')$. The second possibility could lead to significant mixing between the two Z' bosons. Of course, symmetry breaking induced by the VEVs of the scalar components of $\bar{N}(8)$ or $\bar{N}(8')$ would lead to mixing between the corresponding gauginos and the fermionic sterile neutrinos in $\bar{N}(8)$ and $\bar{N}(8')$.

The composite $SU(2)$-singlet N_4 supermultiplet suggested by the strong coupling of the $Sp(2)_B$ group (Sec. VI) could also play a role in the breaking of the extra $U(1)''$s. This state is only relevant to the effective theory below the scale at which $Sp(2)_B$ becomes strong, but that could be anywhere from a few TeV up to very high scales, so it is possible for $B-L$ to be broken along a D-flat direction at a large scale. Of course, other dynamical mechanisms, such as fermion condensates, could possibly be relevant to the breaking of the extra $U(1)''$s.

V. GAUGE COUPLINGS

The gauge couplings are associated with different stacks of branes and do not exhibit conventional gauge unification. Nevertheless, the value of each gauge coupling at the string scale is predicted in terms of a modulus χ and the ratio of the Planck to string scales. The running is strongly affected by the exotic matter and multiple Higgs fields, leading to low values of the MSSM sector couplings at low energy. However, the hidden sector groups are asymptotically free.

The gauge coupling of the gauge field from a stack of $D6$-branes wrapping a three-cycle is given by [15] (the factors of $\sqrt{2\pi}$ were carefully worked out in [8]):

$$g_{YM}^2 M_P^{(4d)} = \sqrt{2\pi} M_s \frac{\sqrt{V_6}}{V_3},\tag{4}$$

where V_3 is the volume of the three-cycle, V_6 is the total internal volume, and M_s is the string scale. The four-dimensional Planck scale is defined as the coefficient of the Einstein term in the low energy effective action:

$$S_{4d}=(M_P^{(4d)})^2\int dx^4\sqrt{g}R+\cdots$$

$$=\frac{1}{16\pi G_N}\int dx^4\sqrt{g}R+\cdots.\tag{5}$$

Since $G_N^{-1/2}=1.22\times10^{19}$ GeV, we have $M_P^{(4d)}=1/4\sqrt{\pi}\times G_N^{-1/2}=1.7\times10^{18}$ GeV. (Some authors define a scale larger by $\sqrt{2}$.) V_3 is given by

$$V_3=(2\pi)^3\prod_{i=1}^{3}\sqrt{n^{i2}(R_1^i)^2+\hat{m}^{i2}(R_2^i)^2},\tag{6}$$

where $R_{1,2}^i$ are the radii of the two dimensions of the ith two-torus, $\hat{m}^i=m^i$ for $i=1,2$, $\hat{m}^3=\tilde{m}^3$, and the wrapping numbers $(n^i,\hat{m}^i)$ are given in Table I. The total internal volume is given by

$$V_6=\frac{(2\pi)^6}{4}\prod_{i=1}^{3}R_1^iR_2^i.\tag{7}$$

The factor of 1/4 comes from the fact that we are orbifolding T^6 by $\mathbf{Z}_2\times\mathbf{Z}_2$, which is an Abelian group of order 4. The Planck scale and Yang-Mills couplings are related to the string coupling g_s by

$$(M_P^{(4d)})^2=\frac{M_s^8V_6}{(2\pi)^7g_s^2}\tag{8}$$

and

$$\frac{1}{g_{YM}^2}=\frac{M_s^3V_3}{(2\pi)^4g_s}.\tag{9}$$

In terms of the complex structure moduli $\chi_i=R_2^i/R_1^i$,

$$g_{YM}^2=\frac{\sqrt{2\pi}M_s}{2M_P^{(4d)}}\frac{\sqrt{\chi_1\chi_2\chi_3}}{\prod_{i=1}^{3}\sqrt{n^{i2}+\hat{m}^{i2}\chi_i^2}}.\tag{10}$$

Supersymmetry requires that $\chi_1:\chi_2:\chi_3=1:3:2$. Therefore

$$g_{YM}^2=\frac{\sqrt{12\pi}M_s}{2M_P^{(4d)}}\cdot\frac{\chi^{3/2}}{\sqrt{[(n^1)^2+(m^1)^2\chi^2][(n^2)^2+9(m^2)^2\chi^2][(n^3)^2+4(\tilde{m}^3)^2\chi^2]}},\tag{11}$$

where $\chi \equiv \chi_1$.

At a scale M below the string scale, the coupling $\alpha_a = g_a^2/4\pi$ of the ath gauge factor is given (at one loop) by

$$\frac{1}{\alpha_a(M)} = \frac{c_a(\chi)}{\alpha_G(\chi)} + b_a t, \tag{12}$$

where

$$\alpha_G(\chi) = \sqrt{\frac{3}{\pi}} \frac{M_s}{4M_P^{(4d)}} \chi^{3/2} \tag{13}$$

and

$$t = \frac{1}{2\pi} \ln \frac{M_s}{M}. \tag{14}$$

For $M = M_Z$ and $M_s \sim M_P^{(4d)}$ one has $t \sim 6$. The c_a and b_a are displayed in Table VI. For the b_a the contributions $b_a(\text{int})$ from the states at the intersecting branes (used in the analysis) and those $b_a(aa)$ from the unwanted aa states are displayed separately. The hypercharge $U(1)$ is in general a linear combination of $U(1)$'s from branes wrapping around different three-cycles:

$$Q_Y = \sum_a d_a Q_a. \tag{15}$$

TABLE VI. Coefficients c_a of $1/\alpha_G$, and β functions b_a for the ath gauge group. The contributions $b_a(\text{int})$ from the intersecting branes (used in the analysis), and the additional contributions $b_a(aa)$ from unwanted aa states not localized at the intersections are both displayed. $Sp(2)_B$ and $Sp(2)_A$ refer to the $Sp(2)$ groups associated with the $B2$ and $A2$ branes, respectively. The gauge coupling for hypercharge is not independent of that for $B-L$ and $Q_8 \pm Q_{8'}$, but is shown for completeness. For comparison, the β functions $b_a(\text{MSSM})$ in the MSSM for $SU(3)$, $SU(2)$, and $1/\alpha_1$ are, respectively, -3, 1, and $33/5$.

Group a	c_a	$b_a(\text{int})$	$b_a(aa)$
$SU(3)$	$1+\chi^2$	-1	9
$SU(2)$	$1+9\chi^2$	18	6
$B-L$	$\frac{10}{9}(1+\chi^2)$	$\frac{64}{3}$	
$Q_8 \pm Q_{8'}$	$12\chi^2$	80	144
$\frac{1}{\alpha_1} = \frac{3}{5}\frac{1}{\alpha_Y}$	$\frac{1}{6}+\frac{59\chi^2}{30}$	$\frac{76}{5}$	$\frac{108}{5}$
$Sp(2)_B$	$6\chi^2$	-4	
$Sp(2)_A$	2	-6	
$Sp(4)$	$2\chi^2$	-5	3

Its coupling α_Y (or the related $\alpha_1 = \sqrt{5/3}\,\alpha_Y$ conventionally used in the MSSM) is not independent of the three other $U(1)$'s, but is included for convenience. It is given by the relation

$$\frac{1}{\alpha_Y} = \sum_a d_a^2 \frac{1}{\alpha_a}. \tag{16}$$

From Eq. (8), along with the requirement $M_s^6 V_6 \gtrsim (2\pi)^6$ (i.e., the compactification radii cannot be smaller than the string scale) and $g_s^2 \lesssim 4\pi$ (perturbative string theory), we expect that $M_P^{(4d)} \gtrsim M_s/\sqrt{8\pi}$. We will generally take $M_P^{(4d)} = M_s$ in our numerical examples. Near unification of, e.g., the $SU(3)$ and $SU(2)$ couplings at M_s prefers small χ, but the overall scale α_G requires that χ cannot be too small. The best results are for $\chi \sim 0.5$. The predicted MSSM gauge couplings at the electroweak scale are presented as a function of χ in Fig. 1. It is seen that the predicted value of α_3 is quite close to the experimental value for $\chi \sim 0.5$ ($1/\alpha_3 \sim 8.5$). However, α is predicted to be much too small, mainly because of the contributions of the exotic states to the running, while $\sin^2 \theta_W$ is predicted too large by a factor ~ 2. (The predicted value of $\sin^2 \theta_W$ at the string scale M_s ranges from 0.78 to 0.73 for χ varying from 0 to ∞, while $1/\alpha_G \sim 12$ for $\chi \sim 0.5$ and $M_P^{(4d)} = M_s$.)

Although the predictions for the MSSM gauge couplings are unrealistic, the quasihidden sector Sp groups are all asymptotically free. Figure 2 displays the scales at which each Sp group becomes strongly coupled as a function of χ. All three groups become strong above the electroweak scale for $M_P^{(4d)} = M_s$ and $\chi \gtrsim 0.4$. For $M_P^{(4d)} < M_s$ the couplings become strong at higher scales [e.g., $Sp(2)_B$ would become strong at around 10^{15} GeV for $M_P^{(4d)} = M_s/3$ and $\chi = 0.5$, with possible implications for neutrino mass, as mentioned in Sec. IV]. The implications will be further discussed in Sec. VI and in [23].

VI. IMPLICATIONS OF A STRONGLY COUPLED SECTOR

We have seen that for sufficiently low values of $M_P^{(4d)}/M_s$ the quasihidden sector groups $Sp(2)_B$, $Sp(2)_A$, and $Sp(4)$ will become strongly coupled above the electroweak scale. This is likely to lead to supersymmetry breaking in the hidden sector, which can be transmitted to the observable sector by supergravity, as well as dilaton/moduli stabilization [23]. Here we will focus on another aspect, i.e., the confinement of free Sp charges, expected in analogy with QCD. In particular, it is plausible to assume that when one of the $Sp(n)$ groups becomes strongly coupled at some scale $M \gg M_Z$ any states carrying $Sp(n)$ charges become confined and drop out of the physical spectrum. However, there may be Sp-neutral bound-state chiral supermultiplets remaining in the spectrum,[4] which may be required to avoid the introduction of anomalies for the remaining gauge groups.

$Sp(4)$ is expected to become strong at a high scale. The

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[4] Supersymmetry breaking associated with a related gaugino condensation will be transmitted to the ordinary sector only by weak supergravity effects, and can be ignored for the purposes of the present discussion.
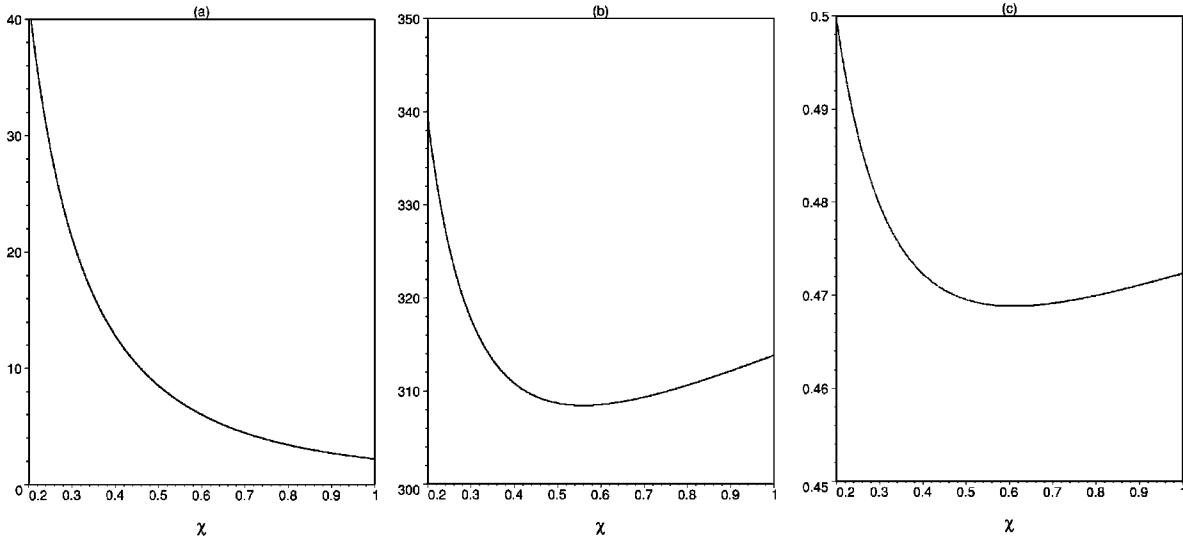
FIG. 1. Predicted values of (a) $1/\alpha_3$, (b) $1/\alpha = 1/\alpha_Y + 1/\alpha_2$, and (c) the weak angle $\sin^2 \theta_W = \alpha_Y/(\alpha_2 + \alpha_Y)$ at the electroweak scale as a function of χ for $M_P^{(4d)} = M_s$. Only the contributions $b_a(\text{int})$ of the states localized at the brane intersections (as well as the gauge bosons) to the running are included. The experimental values are ~ 8.5, 128, and 0.23, respectively.

only state in the chiral sector with $Sp(4)$ charge is from the $B_1 C_2$ intersection. We denote this state as Ψ, which transforms as (1, 2, 1, 1, 4) under $SU(3) \times SU(2) \times Sp(2)_B \times Sp(2)_A \times Sp(4)$. The strong $Sp(4)$ forces may lead to a composite chiral supermultiplet Ψ_{AA}^2, where the subscript indicates an antisymmetrization in both the $SU(2)$ and $Sp(4)$ indices. Ψ_{AA}^2 is a total gauge singlet. Whether or not this composite state is formed, the confinement of Ψ does not lead to any anomalies for the residual gauge groups. An anomaly is induced in the $Q_2 - SU(2)^2$ vertex, where Q_2 is associated with an anomalous $U(1)$ with a massive gauge boson. This can be regarded as a breaking of Q_2 and is presumably harmless, analogous to the breaking of the global axial $U(1)$ symmetry in QCD. The four $SU(2)$ doublets contained in Ψ drop out of the renormalization group equations for $SU(2)$ at the decoupling scale, lowering $b_{SU(2)}$ by 2, but this has only a minor impact on the discussion in Sec. V.

$Sp(2)_B$ can become strongly coupled anywhere from a few TeV up to very high scales such as 10^{15} GeV, depending on χ and $M_P^{(4d)}/M_s$. The fractionally charged $B_2 C_1$ states are charged under $Sp(2)_B$. Let us denote them as $\Phi \equiv (3,1,2,1,1)$ and $\Sigma \equiv (1,1,2,1,1)$. By our assumptions, these will be confined at the $Sp(2)_B$ scale. The strong $Sp(2)_B$ binding might form the composite color triplet $\Phi\Sigma = (3,1,1,1,1)$. This has charge $-1/3$, and would be a candidate for an exotic $SU(2)$-singlet down-type quark, except that it has lepton number $L=1$. Furthermore, if this were the only massless composite, anomalies would be induced in the Q_Y^3 and $Q_Y - SU(3)^2$ vertices. The anomaly-matching condition suggests that, instead of $\Phi\Sigma$, there is a more complicated spectrum of massless composites. The simplest possibility is that the spectrum consists of

$$E_4 = \Phi\Sigma\bar{U},$$

$$N_4 = \Phi\Sigma\bar{D},$$

$$U_4 = \Phi\Sigma\bar{E}, \qquad (17)$$

$$D_4 = \Phi\Sigma\bar{N},$$

i.e., $\Phi\Sigma$ forms bound states with each member of one of the four families of $SU(2)$-singlet antiquarks and antileptons. The latter do not have to drop out of the spectrum, so the
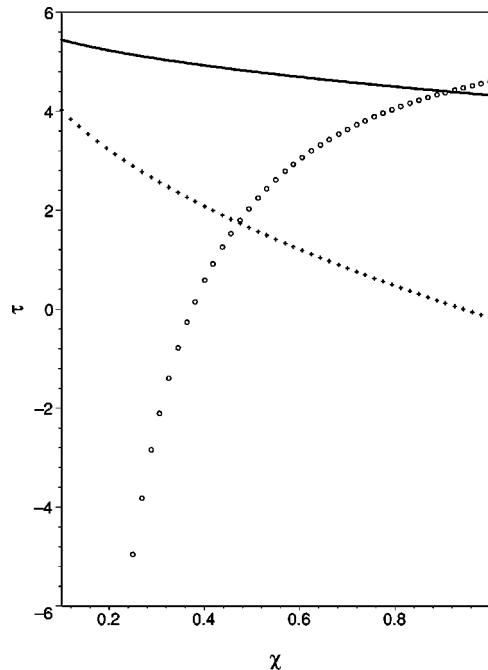

FIG. 2. Scales $\tau = 1/2\pi \ln M/M_Z$, where M is the scale at which the Sp group becomes strongly coupled for $M_P^{(4d)} = M_s$. The curves are $Sp(2)_A$ (solid), $Sp(2)_B$ (circles), and $Sp(4)$ (crosses).

composites U_4, D_4, E_4, and N_4 are candidates to be the exotic [$SU(2)$-singlet] left-handed partners of the elementary fourth family of antiparticles. This spectrum matches all of the gauge anomalies, although it does lead to a (presumably harmless) $(B+L)-Q_8^2$ anomaly, again similar to the axial $U(1)$ in QCD.

This binding mechanism seems very plausible from the viewpoint of anomaly matching. However, it is harder to understand from the actual forces between the constituents, since the $\bar{U}$, $\bar{D}$, $\bar{E}$, $\bar{N}$ are not charged under the strong $Sp(2)_B$ group. (They do carry other gauge charges.)

The decoupling of Φ and Σ and the appearance of the composite states lead to a net increase of 6/5 for the β function for $1/\alpha_1$ at the decoupling scale ($b_{SU(3)}$ and $b_{SU(2)}$ do not change), but this is a small effect for the specific numerical example we have displayed. For example, for $M_P^{(4d)} = M_s$ and $\chi = 0.5$, $Sp(4)$ and $Sp(2)_B$ become strong at $\sim 10^{15}$ and 2×10^6 GeV, respectively. Including both decouplings, the predicted values of α^{-1} and $\sin^2 \theta_W$ decrease by ~ 6 and 0.02, respectively, compared to those in Fig. 2.

$Sp(2)_A$ may also become strongly coupled. However, there are no chiral states with $Sp(2)_A$ charges.

VII. DISCUSSION

In this paper we have described the phenomenological implications of a semirealistic supersymmetric three family model derived from an orientifold construction. In addition to the MSSM, the model involves an extended gauge structure, including two additional $U(1)'$ factors, one of which has family nonuniversal couplings. There is also a quasihidden sector non-Abelian group, which becomes strongly coupled above the electroweak scale. There are many exotic chiral supermultiplets, including an exotic [$SU(2)$-singlet] fourth family of quarks and leptons in which the left-chiral states have unphysical fractional electric charges. These are presumably confined by the strong hidden sector interactions, while anomaly constraints imply composite left-chiral states with the correct charges. The right-chiral states are elementary. The Yukawa sector [22] and other aspects of the hidden sector [23] will be presented separately.

As emphasized in the Introduction, none of the models that have been constructed are fully realistic, and it is difficult to know whether the specific features of a given model are hints of possible new TeV scale physics, or merely artifacts of the construction. For that reason, it is useful to contrast some of the features of this orientifold construction with those of a specific heterotic model described in [4]. For convenience, those predictions are described in more detail in the Appendix.

Both models predict additional $U(1)'$ gauge symmetries, some of which have family-nonuniversal and therefore flavor changing neutral current couplings. Both are most likely broken at the TeV scale, but have a possibility of being broken at an intermediate scale along a D-flat direction. Both also have quasihidden non-Abelian gauge sectors. This means that while most of the states in the model are charged under one or neither of the gauge sectors, there are a few states which couple to both. [The $U(1)'$ also connect the two sec-

tors.] These mixed states have fractional charges like 1/6 or $\pm 1/2$. A hidden sector is an ideal candidate for dynamical supersymmetry breaking if it is strongly coupled. In the heterotic example the hidden sector groups are not asymptotically free. However, in the orientifold example, the groups are asymptotically free, and may lead to gaugino condensation, dilaton/moduli stabilization, and charge confinement, modifying the low energy spectrum.

Both models involve exotic states, often with no satisfactory means of generating fermion masses. These include additional Higgs doublets and singlets, suggesting such effects as a rich spectrum of Higgs particles, neutralinos, and charginos, perhaps with nonstandard couplings due to mixing and flavor changing effects. The effective μ terms are either missing or nonstandard. There may also be vector pairs of additional quarks and leptons. In the orientifold case, the left-chiral states are composite and their right-chiral partners elementary. The orientifold model also contains unwanted adjoints. There may be mixing between lepton and Higgs doublets, leading to lepton number violation. This was required in the heterotic case (where baryon number violation was also possible for one flat direction) and optional for the orientifold.

Although the Yukawa couplings have tree-level contributions in string perturbation theory (in orders of g_s) in the two constructions, they have different origins from the worldsheet perspective. In the orientifold model, the Yukawa couplings arise from nonperturbative effects (worldsheet instantons) in the worldsheet conformal field theory. In the Chaudhuri, Hockney, and Lykken (CHL5) model, the Yukawa couplings are tree-level from the worldsheet point of view. We note, however, that in some heterotic string constructions, the quarks, leptons, and the Higgs fields are localized at different orbifold fixed points (see, e.g., [31] and references therein). In these constructions, the Yukawa couplings also come from worldsheet instantons. Both constructions can yield masses and mixings for some, but not all, of the fermions, but the details depend on the mechanism of supersymmetry breaking. Neither has an obvious mechanism for a neutrino seesaw, except possibly for the case of an intermediate scale breaking of a $U(1)'$.

In the heterotic model the gauge unification predictions are nonstandard (and not very successful) due to the combination of exotic particles contributing to the running of the gauge couplings and higher Kač-Moody levels at the string scale. The orientifold predictions are also nonstandard: the gauge groups are associated with different branes, and have nonstandard moduli-dependent boundary conditions at the string scale, and there are also exotic particles contributing to the running, leading to electroweak couplings that are too small. Resolutions of these difficulties in more successful constructions might involve avoiding these nonstandard features, having exotics which fall into complete grand unification multiplets, or invoking cancellations of effects occurring by accident or due to some other mechanism.

APPENDIX

As described in the Introduction, there has been considerable study of semirealistic perturbative heterotic string constructions, including a class of free-fermionic string models which contain the gauge group and matter content of the MSSM [2,3]. Such constructions generally involve additional gauge factors and many extra matter supermultiplets. However, they also contain an anomalous $U(1)_A$ symmetry and a corresponding Fayet-Iliopoulos contribution to the $U(1)_A$ D term. Maintaining supersymmetry at the string scale requires that some of the fields in the effective four-dimensional theory must acquire compensating vacuum expectation values (VEVs) near the string scale, while maintaining F flatness and D flatness for the other gauge factors. These break some of the extra gauge symmetries, remove some of the apparently massless states from the low energy effective theory, and require that the theory be restabilized, i.e., the superpotential for the remaining massless states must be recalculated when some of the fields are replaced by their string-scale VEVs. A systematic procedure was developed in [32] to classify the flat directions associated with non-Abelian singlet fields. This was used in [4] to investigate the flat directions and related low energy phenomenology in detail for a promising model due originally to Chaudhuri, Hockney, and Lykken [3] (CHL5). Flat directions associated with nonsinglets were studied in [33]. The procedures were used to study the flat directions in a class of models due to Faraggi, Nanopoulos, and Yuan [2] in [34].

The features of perturbative heterotic string models are illustrated by the prototypical example of the CHL5 model. These include the following.

One or more additional (nonanomalous) $U(1)'$ gauge factors. The associated Z' gauge boson is typically expected to be lighter than around 1 TeV [27], although for one flat direction studied the breaking could be at an intermediate scale [30]. The Z' couplings are family nonuniversal, leading to flavor changing neutral currents (FCNC) [28].

Additional non-Abelian gauge factors. These could, in principle, play a role in dynamical supersymmetry breaking. However, in the model studied the factors do not become strongly coupled below the string scale (i.e., they are not asymptotically free). These extra gauge factors are quasihidden, i.e., most matter multiplets transform nontrivially under the standard model $SU(3) \times SU(2)$ group or under the hidden sector group (or neither), but not both. However, there are a few exceptions which are charged under both. The extra $U(1)'$ typically couple to both sectors.

There are many exotic supermultiplets in the model, including an extra d-type quark, extra Higgs/lepton doublets, and many non-Abelian singlets. For many, there is no mechanism to give a significant mass to the fermions. This is a major flaw of most such models. (One exception, which, however, has unrealistic Yukawa couplings, is described in [34].) The spectrum includes a number of charge $\pm 1/2$ states. These are all charged under the quasihidden sector group, and potentially could disappear from the spectrum if the hidden sector charges are confined [35]. However, as noted above, the hidden sector factors in the CHL5 model are not strongly coupled, so this mechanism does not occur.

There are more than the two Higgs doublets of the MSSM. In addition to the more complicated Higgs spectrum, there is a possibility of Higgs mediated FCNC. Models with an electroweak scale $U(1)'$ can generically provide a natural solution to the μ problem [26,27], which is generated dynamically by the VEV of the field that breaks the $U(1)'$. However, in the CHL5 model the effective μ terms are noncanonical, connecting Higgs doublets which generate the t and b masses to others. In the specific examples studied, one of the needed effective μ terms is absent, leading to an unwanted global symmetry.

The model has gauge coupling unification. However, the detailed predictions for the low energy couplings differ from the MSSM (and from experiment) because of the additional matter fields as well as higher Kač-Moody levels for the $U(1)$ factors.

The Yukawa couplings at the string scale are either g, $g/\sqrt{2}$, or 0, where $g = O(1)$ is the gauge coupling, allowing large masses for the t and b, and the possibility of radiative electroweak breaking. Smaller Yukawa couplings can be associated with higher dimensional operators that become cubic after vacuum restabilization. CHL5 contains $t-b$ and an unphysical $\mu - \tau$ universality, a noncanonical $b-\tau$ relation, and a nontrivial (but unphysical) d quark texture for one flat direction.

The models violate R parity and lepton number (due to mixing between lepton and Higgs doublets), so there is no stable neutralino. The Baryon number is violated for one flat direction, leading to proton decay and $n-\bar{n}$ oscillations (with rates that cannot be calculated without resolving the problem of the massless exotics).

There is no obvious mechanism for a neutrino seesaw except for the cases in which one of the $U(1)'$ gauge factors is broken at an intermediate scale [30].

When phenomenological soft supersymmetry breaking parameters are introduced by hand at the string scale, one can calculate the symmetry breaking and the spectrum of the Higgs and supersymmetry particles. One can obtain a sufficiently large Z' mass (e.g., >700 GeV) and small $Z-Z'$ mixing (e.g., <0.005) for somewhat tuned values of the soft parameters. The large Z' mass scale is set by the soft breaking parameters, implying a spectrum quite different from most of the parameter space of the MSSM: typically, the squark and slepton masses are in the TeV range (except possibly for the third family), but there is a richer spectrum of Higgs particles, charginos, and neutralinos [36].

[1] For reviews, see, e.g., B. R. Greene, in *High Energy Physics and Cosmology*, Proceedings of the 1990 Summer School, Trieste, Italy, 1990, edited by J. C. Pati *et al.* (World Scientific, Singapore, 1991); F. Quevedo, hep-th/9603074; A. E. Faraggi, hep-ph/9707311; Z. Kakushadze, G. Shiu, S. H. Tye, and Y. Vtorov-Karevsky, Int. J. Mod. Phys. A **13**, 2551 (1998); G. Cleaver, M. Cvetič, J. R. Espinosa, L. L. Everett, P. Langacker, and J. Wang, Phys. Rev. D **59**, 055005 (1999), and references therein.

[2] A. E. Faraggi, D. V. Nanopoulos, and K. Yuan, Nucl. Phys. **B335**, 347 (1990).

[3] S. Chaudhuri, G. Hockney, and J. Lykken, Nucl. Phys. **B469**, 357 (1996).

[4] Cleaver, Cvetič, Espinosa, Everett, Langacker, and Wang [1]; Phys. Rev. D **59**, 115003 (1999).

[5] C. Angelantonj, M. Bianchi, G. Pradisi, A. Sagnotti, and Ya. S. Stanev, Phys. Lett. B **385**, 96 (1996).

[6] M. Berkooz and R. G. Leigh, Nucl. Phys. **B483**, 187 (1997).

[7] Z. Kakushadze and G. Shiu, Phys. Rev. D **56**, 3686 (1997); Nucl. Phys. **B520**, 75 (1998); Z. Kakushadze, *ibid.* **B512**, 221 (1998).

[8] G. Shiu and S.-H. H. Tye, Phys. Rev. D **58**, 106007 (1998).

[9] Z. Kakushadze, G. Shiu, and S. H. Tye, Nucl. Phys. **B533**, 25 (1998); Phys. Rev. D **58**, 086001 (1998).

[10] G. Aldazabal, A. Font, L. E. Ibáñez, and G. Violero, Nucl. Phys. **B536**, 29 (1999).

[11] Z. Kakushadze, Phys. Lett. B **434**, 269 (1998); Phys. Rev. D **58**, 101901 (1998); Nucl. Phys. **B535**, 311 (1998).

[12] M. Cvetič, M. Plümacher, and J. Wang, J. High Energy Phys. **04**, 004 (2000).

[13] M. Klein and R. Rabadán, J. High Energy Phys. **10**, 049 (2000).

[14] M. Cvetič, A. M. Uranga, and J. Wang, Nucl. Phys. **B595**, 63 (2001).

[15] M. Cvetič, G. Shiu, and A. M. Uranga, Phys. Rev. Lett. **87**, 201801 (2001); Nucl. Phys. **B615**, 3 (2001).

[16] M. Cvetič, G. Shiu, and A. M. Uranga, "Chiral type II orientifold constructions as M theory on G_2 holonomy spaces," hep-th/0111179.

[17] M. Berkooz, M. R. Douglas, and R. G. Leigh, Nucl. Phys. **B480**, 265 (1996).

[18] R. Blumenhagen, L. Goerlich, B. Kors, and D. Lust, J. High Energy Phys. **10**, 006 (2000).

[19] G. Aldazabal, S. Franco, L. E. Ibáñez, R. Rabadán, and A. M. Uranga, J. Math. Phys. **42**, 3103 (2001); J. High Energy Phys. **02**, 047 (2001).

[20] R. Blumenhagen, B. Körs, and D. Lüst, J. High Energy Phys. **02**, 030 (2001).

[21] C. Angelantonj, I. Antoniadis, E. Dudas, and A. Sagnotti, Phys. Lett. B **489**, 223 (2000).

[22] M. Cvetič, G. Shiu, and P. Langacker, UPR-0991T.

[23] M. Cvetič, P. Langacker, and J. Wang (work in progress).

[24] C. Angelantonj and R. Blumenhagen, Phys. Lett. B **473**, 86 (2000).

[25] L. E. Ibáñez, F. Marchesano, and R. Rabadán, J. High Energy Phys. **11**, 002 (2001).

[26] D. Suematsu and Y. Yamagishi, Int. J. Mod. Phys. A **10**, 4521 (1995).

[27] M. Cvetič, D. A. Demir, J. R. Espinosa, L. L. Everett, and P. Langacker, Phys. Rev. D **56**, 2861 (1997); **58**, 119905(E) (1997).

[28] P. Langacker and M. Plümacher, Phys. Rev. D **62**, 013006 (2000).

[29] See, J. Erler and P. Langacker, Phys. Rev. Lett. **84**, 212 (2000), and references therein.

[30] G. Cleaver, M. Cvetič, J. R. Espinosa, L. L. Everett, and P. Langacker, Phys. Rev. D **57**, 2701 (1998); P. Langacker, *ibid.* **58**, 093017 (1998).

[31] M. Cvetič, Phys. Rev. Lett. **59**, 1795 (1987).

[32] G. Cleaver, M. Cvetič, J. R. Espinosa, L. L. Everett, and P. Langacker, Nucl. Phys. **B525**, 3 (1998); **B545**, 47 (1999).

[33] M. Cvetič, L. Everett, P. Langacker, and J. Wang, J. High Energy Phys. **04**, 020 (1999).

[34] G. B. Cleaver, A. E. Faraggi, D. V. Nanopoulos, and J. W. Walker, Nucl. Phys. **B593**, 471 (2001); **B620**, 259 (2002).

[35] See, C. Coriano, A. E. Faraggi, and M. Plümacher, Nucl. Phys. **B614**, 233 (2001), and references therein.

[36] L. L. Everett, P. Langacker, M. Plümacher, and J. Wang, Phys. Lett. B **477**, 233 (2000).