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D-branes and textures

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Abstract: We examine the flavor structure of the trilinear superpotential couplings which can result from embedding the Standard Model within D-brane sectors in type-IIB orientifold models, which are examples within the type-I string framework. We find in general that the allowed flavor structures of the Yukawa coupling matrices to leading order are given by basic variations on the "democratic" texture ansatz. In certain interesting cases, the Yukawa couplings have a novel structure in which a single right-handed fermion couples democratically at leading order to three left-handed fermions. We discuss the viability of such a "single right-handed democracy" in detail; remarkably, even though there are large mixing angles in the u, d sectors separately, the CKM mixing angles are small. The analysis demonstrates the ways in which the type-I superstring framework can provide a rich setting for investigating novel resolutions to the flavor puzzle.

Keywords: D-branes, Superstring Vacua, Supersymmetric Effective Theories, Quark Masses and SM Parameters.

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1. Introduction

Uncovering the nature and origin of the fermion mass hierarchy and mixings is one of the most fundamental issues in high energy physics. This topic has thus been the subject of intensive research effort, which has yielded a wealth of phenomenological literature exploring the possible textures of the quark Yukawa coupling matrices consistent with the experimentally determined quark masses and Cabibbo-Kobayashi-Maskawa (CKM) mixings [1]. The recent SUPER-KAMIOKANDE results [2], which have provided overwhelming evidence for atmospheric neutrino oscillations, have also opened the door for explorations of possible textures in the lepton sector [3]; as is well known, the fact that maximal mixing is favored between the second and third generations of the lepton sector leads to very different possibilities for textures than in the quark sector (where all mixing angles are small).

Despite the insight which can be gained from these phenomenological studies of the fermion mass matrices, arguably the true resolution to the flavor problem lies in the domain of the fundamental theory. Since at present superstring/"M" theory is the only candidate for a truly fundamental quantum theory of all interactions, studies of the flavor structure of the Yukawa couplings within four-dimensional superstring models are well motivated. In addition, the couplings of the effective lagrangian in superstring theory are *calculable* (at least in principle), and *not input* parameters. This important feature allows for the flavor problem to be addressed quantitatively within a given superstring model, without ad hoc assumptions or the necessary introduction of small but arbitrary parameters.

In perturbative heterotic string models, the flavor structure of the Yukawa couplings has been studied extensively in both explicit quasi-realistic 4D constructions such as orbifold models [4] and free fermionic models [5, 6, 7] (with some but limited success due to the absence of fully realistic models), as well as in string-motivated effective QFT's [8]. An important result of the analyses of explicit string models was to demonstrate that the trilinear superpotential couplings at the string scale are generally either zero or $\mathcal{O}(1)$, such that they can provide a natural explanation for the top quark Yukawa coupling [6]. Several mechanisms utilizing higher-dimensional operators (which again are calculable in string theory) are then available for the generation of the lighter Yukawa couplings,[1] and have been explored both within the string-derived and string-motivated approaches. Perhaps the most popular approach has been to utilize the anomalous U(1) gauge symmetry and the associated vacuum stabilization procedure generic to perturbative heterotic models [4, 6, 8, 10].

Such perturbative heterotic models have traditionally been thought to be strong candidates for realistic superstring compactifications and thus issues such as the flavor problem have been explored most extensively in this context. However, newer classes of perturbative string vacua which have potentially very different phenomenological properties have recently become accessible for study with the advent of duality symmetries and the discovery of Dirichlet branes. We focus here on the $N = 1$, $D = 4$ type-IIB orientifold models [12]–[18], which are the simplest examples within the more general type-I string picture. In these models, the gauge groups of the effective low-energy lagrangian arise from sets of coincident D-branes and the matter fields (such as the MSSM fields) arise from open strings which must start and end upon D-branes, such that the phenomenological implications depend quite crucially on the nature of the embedding of the SM gauge groups within the different D brane sectors. We will see this feature will provide a new framework to address issues of flavor physics and can lead to novel Yukawa textures dictated in part by the nature of the SM gauge group embedding.

We emphasize that despite significant recent advances (see e.g. [16]), the development of model-building techniques for type-I models is still at early stages, and there are but few quasi-realistic models (see, however, [14]–[17]). In contrast, investigations into the general phenomenological features of these models have been possible recently within the string-motivated approach, which is the approach we adopt in this paper. This advance has been due to Ibáñez, Muñoz, and Rigolin [18], in which a classification of the matter fields and the structure of the tree-level couplings of the effective lagrangian of these fields have been extracted on general grounds. The authors of [18] further write down the form of the soft supersymmetry breaking mass

[1] In orbifold models, the trilinear couplings of $\mathcal{O}(1)$ are those among untwisted sector states. For couplings involving twisted sector fields, the Yukawa couplings can have additional suppressions (see e.g. [4]), such that in principle it might be possible to obtain the appropriate hierarchies from the original trilinear couplings only; see e.g. [9] for an analysis within the string-motivated approach.

parameters (assuming the dilaton/moduli play the dominant role in SUSY breaking), which has enabled a number of studies of the patterns of soft breaking parameters in this class of models [18, 19, 20]. Such models are very interesting, in that they provide an attractive alternative to older approaches to collider and dark matter phenomenology, and also can address CP violation issues.

In this paper, we present a study of Yukawa textures arising from D brane theories, and discuss the implications for the soft mass parameters. Throughout this paper, we assume all extra dimensions are small, and that the string and unification scales coincide. We are mainly concerned with the large (order one) Yukawa couplings in these theories, although we give a brief discussion of possible theoretical origins of small Yukawa couplings below. Our main observation is that in these theories, the typical leading order Yukawa couplings do not consist of the traditional hierarchical textures with a single entry in the 33 position, but instead tend to predict textures which at lowest order are either democratic or involve a single right-handed fermion coupling democratically to the three left-handed families. Of course, after a field rotation, such structures at leading order are equivalent to the hierarchical structure, but from the point of view of calculating sub-leading perturbations, it is most natural to work in the original basis defined by the theory. Furthermore, the soft mass matrices are also determined in the basis determined by the theory, and if they have large off-diagonal elements in the original basis, then they will retain large off-diagonal elements in the hierarchical basis, with important implications for flavor physics. Thus, the democratic and hierarchical bases are not equivalent but are physically distinguishable.

For the purposes of this study, what is most important is the form of the trilinear (Yukawa) superpotential couplings, which are known [13, 18] and have a natural interpretation in terms of the interactions of open strings with D-branes [13]. As is typical within superstring models, string symmetries can forbid gauge-allowed terms, in contrast to the case within ordinary four-dimensional QFT's. The actual Yukawa couplings from the trilinear superpotential terms (calculated at tree-level in the string loop expansion) of the effective low energy theory are numbers of $\mathcal{O}(1)$ (much like in the untwisted sector of the heterotic orbifold models [18]). In addition, there are no higher string loop (genus) corrections to the superpotential due to supersymmetric nonrenormalization theorems.[2] Hence, for a given embedding of the SM gauge group it is possible to learn about the *lowest-order structure* of the Yukawa coupling matrices in flavor space.

We now briefly discuss the theoretical origin of the small Yukawa couplings in these models. As in the case of the majority of the perturbative heterotic models, the small Yukawa couplings must be obtained from higher-dimensional operators.

[2] See e.g. [11] for the argument within the heterotic case. The same argument largely holds in the type-I theory, since both moduli controlling the gauge couplings of the nine-branes and five-branes are proportional to the genus-expansion parameter.

In general texture schemes, certain couplings must be allowed and others forbidden based on the leading-order form of the Yukawa matrices; this can occur within superstring models either by string selection rules or by gauge invariance (e.g., the charge assignments of the fields with respect to the anomalous U(1)'s can conspire to forbid couplings otherwise allowed by string selection rules). However, the allowed structure of the higher-dimensional operators has not yet been fully explored in the type-I framework, in contrast to the case within perturbative heterotic models. Furthermore, although it is known that the phenomenological implications of the multiple anomalous U(1)'s generic to the type-IIB orientifold models are likely to differ significantly with the situation in perturbative heterotic models,[3] detailed phenomenological studies have not been performed (see however [18, 21, 22, 23]), due in part to the paucity of quasi-realistic models. Therefore, it is difficult to make strong statements about the origin of the higher-order corrections to the Yukawa matrices at this stage of the investigation. Later on, we will present the phenomenological requirements on the form of these corrections, and discuss the possibilities for obtaining such corrections within the type-I framework.

We consider possible Yukawa textures which can arise in such type-I/D-brane models in which the Standard Model gauge group is split between two D-brane sectors. Such scenarios represent simple possibilities for model building, and have inspired previous interest in the realm of supersymmetric CP violation, since it has been shown [19] that certain embeddings can provide for interesting models of the soft breaking parameters which can have large CP-violating phases with small electric dipole moments (EDM's) due to the presence of nontrivial relative phases in the gaugino mass parameters. In fact, the question of the flavor structure of the soft breaking parameters such models has also quite recently been investigated within the context of effects of such large phases on FCNC and CP violation in the neutral meson systems [27]. We emphasize that (putting aside the issue of CP violation) such scenarios also provide an intuitive framework for studying flavor physics, as the different generations can be distinguished in a stringy manner by their quantum numbers with respect to the various D-brane sectors.

[3] As is well known, the situation with the anomalous U(1)'s in the type-I framework differs in several ways from that of the single anomalous U(1) in perturbative heterotic models. In the heterotic case, the triangle anomalies due to the single anomalous U(1) are cancelled by the universal Green-Schwarz (GS) mechanism, in which the dilaton superfield shifts under the anomalous U(1). Since the dilaton is forced to have a nonzero VEV in perturbative heterotic string theory (this VEV gives the gauge coupling), a nonzero Fayet-Iliopoulos term is generated at the string one-loop order; the presence of the FI term triggers certain scalar matter fields to acquire VEV's of $\mathcal{O}(M_{\rm String})$, leading to a supersymmetric ("restabilized") string vacuum. In the type-IIB orientifold models, there are multiple anomalous U(1)'s, which are cancelled not by the universal GS mechanism, but rather by shifts of certain (twisted sector) moduli fields. Unlike the dilaton, these fields are not required to have nonzero VEV's (nonzero VEV's imply a smoothing out ("blowing-up") of the orbifold singularities); thus, nonzero FI terms and the corresponding vacuum restabilization procedure are not necessarily generated due to anomaly cancellation.

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In this paper, we will consider various possibilities for such embeddings, focusing first on textures in the quark sector. We will show that the leading-order flavor structure of the Yukawa matrices $Y_{u,d}$ are in general given by variations on the ansatz of democratic textures [26], in which each nonzero entry of the matrices is accurately equal to unity, times an overall factor $\mathcal{O}(1)$. We will start with an analysis of several representative models in which SU(2) and SU(3) arise from different D-brane sectors, which have been discussed in [19]. We will focus in particular on the case (first written in [19]) in which it is assumed that SU(3) and $\mathrm{U}(1)_Y$ arise from a single D-brane sector, while SU(2) arises from a different D-brane sector. The allowed quark Yukawa couplings in this interesting case can naturally have a novel structure at leading order in which the ratios of the 13, 23, and 33 elements are all accurately equal to unity, such that the up-quark matrix at lowest order is

$$Y_u = \mathcal{O}(1) \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 1 \\ 0 & 0 & 1 \end{pmatrix}, \tag{1.1}$$

where the equal entries in the third column corresponds to the third family right-handed quark coupling democratically to the three left-handed quark families. We then turn to an analysis of this new scenario, which we call the "single right-handed democratic" texture. Note that the product $Y_u Y_u^\dagger$ yields the democratic texture ansatz. The down-quark matrix Y_d has a similar structure at lowest order except that the overall factor may be suppressed for low values of $\tan\beta$, the ratio of Higgs vacuum expectation values. The model yields a new mechanism for the generation of the small CKM angles, which arise from an accurate cancellation of large rotation angles between the up and down sector; as we shall see, this does not involve fine-tuning since the leading order Yukawa couplings are exactly equal at lowest order (string tree-level).

We will also comment on the situation for the lepton sector within this class of D-brane models. In this case the phenomenological requirements (such as large 23 mixing) place a different set of requirements on the form of the Yukawa couplings of the charged leptons and the right-handed neutrinos, as well as on the possible Majorana mass terms for the right-handed neutrinos (needed for the implementation of the see-saw mechanism). The lepton sector will prove to be more challenging, in part because the D-brane assignments of the right-handed neutrinos (which are SM gauge singlets) are not dictated by the SM gauge group embedding, in contrast to the case of the MSSM fields. In the D-brane model described above which leads to the single right-handed democracy for the quarks, we find that if the D-brane assignments for the charged lepton singlets and the right-handed neutrinos mirror those of the quark singlets, one is led to "single right-handed neutrino dominance" (SRHND), as suggested by one of us on purely phenomenological grounds in [28]. However, unlike the case in [28] (in which the numerical values of the couplings of the dominant

RH neutrino to the three lepton doublets were allowed to differ substantially), the constraint of large 23 mixing means in this case that the natural cancellation mechanism referred to above must be thwarted somehow (which appears to be quite a difficult task). We presume this issue is closely related to the different ways that the right-handed quarks and right-handed neutrinos are necessarily represented in the type-I/D-brane theory. In fact, viable lepton textures seem to require alternate SM gauge group embeddings than those of the models considered in this paper, and/or additional suppression mechanisms (e.g., from U(1) family symmetries, etc.). We will present a more complicated D-brane model in which D-brane assignments distinguish the lepton and Higgs doublets which can provide an example where large 23 mixing angles can be obtained without invoking any additional family symmetries, and comment on the implications for further work.

2. Theoretical framework

For completeness of presentation, we first turn to a brief review of the properties of the $D = 4$, $N = 1$ type-IIB orientifold models (but refer the reader to [18] and references therein for a more comprehensive discussion). This class of models consists of orientifold compactifications of the Type IIB theory [12, 14, 15, 18]. Although the starting point is the type-IIB theory of closed superstrings, consistency conditions (tadpole cancellation) require the addition of open string (type-I) sectors and D-branes, upon which the open strings must end. The number and type of D-branes required in a given model will depend on the particular orientifold group; however, in the most general situation there are one set of nine-branes and three sets of five-branes (5_i), in which the index i labels the complex coordinate of the internal space included in the world-volume of the five-brane. We consider this most general case, and further assume that all five-branes are located at a single orbifold fixed point, which leads to enhanced gauge symmetries and a manifestly T-dual spectrum (see e.g. [13, 18]). In these models, each set of coincident D-branes gives rise to a (generically non-abelian) gauge group, such that the SM gauge group is to be embedded in the generic gauge group structure $\mathcal{G} = \mathcal{G}_9 \times \Pi_i \mathcal{G}_{5_i}$.

The chiral matter fields (from which the MSSM matter fields will be obtained) and gauge bosons arise from the open string sectors, and can be classified into two basic categories. The first category consists of open strings which start and end on D-branes of the same sector, for which the corresponding matter fields, which are denoted (in the notation of [18]) as C^9_j, $C^{5_i}_j$ (with $i, j = 1, \ldots, 3$) for the nine-branes and five-branes respectively. These states are charged under the gauge group of the single set of D-branes (typically in the fundamental or antisymmetric tensor representations). The matter fields in the second category, denoted by C^{95_i}, $C^{5_i5_j}$, consist of open strings which start and end on different sets of branes. These fields can thus have quantum numbers (typically in bifundamental representations) with respect to the gauge groups of both D-brane sectors. It is clear that the identification of the MSSM fields from the states of these two categories (e.g., whether the up-type quark singlets $U^c_{1,2,3}$ are $C_1^{5_1}$ states or $C^{5_15_2}$ states, etc.) will depend quite crucially on the nature of the embedding of the SM gauge group within the gauge groups of the D-brane sectors.

The chiral matter fields C_j^9, $C_j^{5_i}$ associated with open strings which start and end on D-branes of the same sector carry an additional index j, which labels the three complex compact dimensions. This additional label can play an important role in the analysis of the Yukawa couplings because the fields which differ only in this index have different couplings in the effective lagrangian [18]. Thus, this additional index can provide for a string suppression of otherwise gauge-allowed terms in certain cases (such as in models based on the Z_3 orbifold in which such fields have identical gauge quantum numbers). The structure of the general trilinear superpotential couplings in this class of models demonstrates this feature explicitly, as can been seen from the expression for the superpotential in [18]. To emphasize this point, we present the relevant superpotential terms for the case of interest in which the SM group will be embedded within at most *two* different D-brane sectors:

$$\begin{aligned} W = & C_1^9C_2^9C_3^9 + \sum_{i=1}^{3} C_1^{5_i}C_2^{5_i}C_3^{5_i} + \sum_{i=1}^{3} C_i^9 C^{95_i}C^{95_i} + \\ & + \sum_{i=1}^{3} C_i^{5_i}C^{95_i}C^{95_i} + \sum_{i\neq j\neq k}^{3} C_k^{5_i}C^{5_i5_j}C^{5_i5_j} \, . \end{aligned} \tag{2.1}$$

The Yukawa couplings (which are not displayed explicitly) involve numerical factors (including traces over the Chan-Paton indices) and are in general numbers of $\mathcal{O}(1)$. In addition, the Yukawa couplings are rescaled upon passing to the low energy theory after SUSY breaking and normalizing the matter fields to have canonical kinetic terms (see e.g. [24]), such that the Yukawa couplings of the low energy theory also depend on the gauge coupling of the relevant D-brane sector gauge group (see e.g. the notation in [18]). It is important to note that the Yukawa couplings of the low energy theory ultimately depend (via this rescaling) on the full Kähler potential of both the moduli and matter fields. Although the tree-level form of the Kähler potential is known (see [18] for the explicit expression), the Kähler potential is not protected by supersymmetric nonrenormalization theorems and gets corrected order by order in string (genus) perturbation theory, such that the detailed numerical values of the Yukawa couplings are not fully under theoretical control. However, the numerical details of the Yukawa couplings are not crucial for the purposes of this study. Rather, what will be important is the fact that at the string tree level, the couplings of the matter fields only distinguish between fields with different D-brane assignments, and do not contain any further information about flavor. Whether or not this feature can be preserved in higher-genus corrections to the Kähler potential is left to a future study.

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Although the relevant couplings have been presented above both for the case in which the SM group is split between two five-brane sectors as well as the case in which the SM group is split between the nine-brane and one five-brane sector, we emphasize that from the phenomenological point of view in general (and certainly for the purposes of this or any similar study) these cases are completely equivalent. This is clear from the manifest T-duality between the different D-brane sectors for the case we study here as pointed out in [18]; even if T duality is spontaneously broken, it is always possible to find a vacuum in the T-dual picture with identical phenomenology. Therefore, in what follows we shall take the case of the 5_1 and 5_2 brane sectors for the sake of definiteness, keeping in mind that these results can be interpreted straightforwardly in T-dual pictures in which the various D-brane sectors are interchanged (we will present explicit examples below).

The superpotential couplings in this case to be investigated are then given by

$$W = \sum_{i=1}^{2} C_1^{5_i} C_2^{5_i} C_3^{5_i} + \sum_{i=1}^{2} C_3^{5_i} C^{5_1 5_2} C^{5_1 5_2} \,. \tag{2.2}$$

Note that eq. (2.2) demonstrates the absence of potentially gauge-allowed terms (such as $C_{1,2}^{5_1} C^{5_1 5_2} C^{5_1 5_2}$, or $C_1^{5_1} C_1^{5_1} C_2^{5_1}$, etc.). In fact, the form of the superpotential is quite restrictive (since it contains only the four types of couplings given above), which will have a significant impact on the texture analysis.

To summarize, the trilinear superpotential couplings given in eq. (2.2) will be the starting point in our analysis. In the upcoming sections, we will investigate different patterns for the Yukawa coupling matrices dictated by the assignments of the MSSM fields to the various D-brane sectors, using the nature of the SM gauge group embedding in each case as our guide. Of course, as the numerical values of the Yukawa couplings in eq. (2.2) are $\mathcal{O}(1)$, the smaller Yukawa couplings must be determined via higher-dimensional operators, either at the string scale or at lower scales in the effective field theory.[4] The allowed higher-dimensional operators present at the string scale (from integrating out heavy string states) certainly also have the feature that certain gauge-allowed combinations are forbidden by string symmetries; hence, it may be possible within our framework of using D-brane assignments to understand which types of operators will lead to the small Yukawa couplings. However, as stated previously a systematic analysis of such operators has not yet been carried

[4]The nature of the couplings of the matter fields in the type-IIB orientifold models is analogous to the case within perturbative heterotic orbifold models in which the MSSM fields are associated with untwisted sector states (states with modular weights $n_i = -1$). However, in perturbative heterotic orbifold models, the MSSM fields can be associated with "twisted" sector states; these states have different modular weights than those of the untwisted sector. In this case, the moduli dependence of the twisted sector Yukawa couplings can allow for the generation of realistic fermion mass matrices at the trilinear order, in sharp contrast to the situation in the type-I models studied in the present paper.

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out (although it is certainly worthy of future investigation). Therefore, we will not be able to say much about the small perturbations to the leading order Yukawas at this time, but instead adopt the strategy of letting phenomenology dictate what small perturbations are required for viable textures.

3. Quark textures and single right-handed democracy

We now consider several simple possibilities for embedding the SM gauge group within the D-brane sectors. Certainly the simplest possibility is to consider that the SM gauge group arises from a single D-brane sector, such that the MSSM fields are most naturally interpreted as the massless open string states which start and end on that set of D branes. In fact, explicit quasi-realistic 4D orientifold models have been constructed recently which have this feature [16]. However, the analysis of the structure of the Yukawa couplings in this case is quite reminiscent of the prototypical cases studied within perturbative heterotic string theory. The reason for this fact is that since all of the MSSM states essentially "belong" to a single D-brane sector, there is no flexibility within the model to distinguish the fields of different families by virtue of their D-brane assignments. Hence, it is necessary to imagine other methods to distinguish between the families which also occur within perturbative heterotic models, such as via different quantum numbers with respect to the anomalous U(1)'s, or other "stringy" quantum numbers which are not inextricably linked to the presence of the D branes.

Since the purpose of this paper is to investigate new possibilities for flavor physics within the type-I string picture, we choose to consider the next simplest case in which the SM gauge group is split between two different D-brane sectors. Within this framework we analyze the case in which the gauge groups SU(3) and SU(2) originate from different five-brane sectors, and consider the two simplest options for $U(1)_Y$:

1. $U(1)_Y$ and SU(2) arise from the same set of branes,
2. $U(1)_Y$ and SU(3) arise from the same D-brane sector.

Such scenarios have nonuniversal gaugino mass parameters with nontrivial relative CP-violating phases (if the fields which break SUSY are assumed to be complex), and therefore have generated interest in the context of obtaining models of the soft SUSY breaking parameters which can have large CP-violating phases but small electric dipole moments of electron and neutron due to cancellations [19]. In particular, it was demonstrated [19] that in the model in which SU(3) and $U(1)_Y$ originate from the same D-brane sector, the EDM's could be small due to cancellations even with $\mathcal{O}(1)$ phases over significant regions of parameter space, while large phases were disfavored in the model with SU(2) and $U(1)_Y$ from the same set of branes.

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Although these two models exhibit very different implications for CP violation, both models have equal interest in the context of flavor physics, in that each dictates interesting leading-order Yukawa coupling matrices which follow from the ways the SM gauge group is embedded and the D brane assignments of the MSSM fields, and thus we will explore each in detail. It is important to note that in contrast to the case in which the SM group is associated with a single set of branes, it is not known whether there are any explicit quasi-realistic orientifold models which display this pattern for the SM gauge group. For example, the SM gauge group is split between two D-brane sectors in the Z_6 model of Shiu and Tye [15], but not in the way given by either of our models; even though SU(3) arises from a single D-brane sector, both SU(2) and $\mathrm{U}(1)_Y$ are given by linear combinations of the (broken) gauge groups of both D-brane sectors. Another important question in this context is the origin of the three families in such configurations with the SM gauge groups arising from different branes; i.e., three copies of (say) $C^{5_1 5_2}$ fields with the same SM gauge quantum numbers may be difficult to obtain in explicit orientifold compactifications. We emphasize that we choose to study these models assuming the presence of three families to simply illustrate issues of flavor physics within the type-I string picture, not because they are candidates for a fully realistic theory. Similar analyses can be done for more complicated SM gauge group embeddings, and we give an example of this later in the discussion of the lepton textures.

We first consider model (1) in which SU(2) and $\mathrm{U}(1)_Y$ originate from the 5_1 sector, while SU(3) is from the 5_2 sector Within this model, this gauge group embedding and the subsequent required D-brane assignments of the MSSM fields unambiguously dictate the form of the Yukawa coupling matrices in the quark sector at leading order. Explicitly, note in this case that *all* states with SU(3) quantum numbers Q_a, U^c_a, D^c_a (in which a is a family index) must be associated with string states from the intersection of the two sets of five-branes; hence the quark doublets and singlets all must be of the form $C^{5_1 5_2}$. The Higgs doublets $H_{u,d}$ must then be states of the type $C^{5_1}_3$ to obtain large top (and bottom) Yukawa couplings, as can be seen from the form of the superpotential in eq. (2.2):

$$W_{\mathrm{quark}} = \sum_{a,b}^{3} [Y^u_{ab} Q_a U^c_b H_u + Y^d_{ab} Q_a D^c_b H_d] \,. \tag{3.1}$$

The couplings of all three families then arise from the $C^{5_1}_3 C^{5_1 5_2} C^{5_1 5_2}$ coupling, such that (in the rescaled Yukawa lagrangian)

$$Y^{u,d} = \mathcal{O}(g_{5_1}) \begin{pmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{pmatrix} , \tag{3.2}$$

which is the "democratic" texture ansatz first explored by Fritzsch and collaborators [26]. As can be seen from the form of eq. (2.1), identical Yukawa matrices occur in the T-dual picture in which SU(2), $\mathrm{U}(1)_Y$ are associated with the nine-brane sector, while SU(3) originates from the e.g. 5_1 sector. The corresponding D-brane assignments are Q_a, U_a^c, D_a^c are C^{95_1} fields, while $H_{u,d}$ are C_1^9 states.

Such a structure can also emerge quite naturally for the quark sector in model (2), in which SU(3) and $\mathrm{U}(1)_Y$ originate from the 5_1 sector, while SU(2) is from the 5_2 sector (first studied in [19]). In this case, all states with SU(2) quantum numbers must be associated with string states from the intersection of the two sets of five-branes; hence the quark doublets Q_a of all three families and the Higgs doublets $H_{u,d}$ all must be of the form $C^{5_15_2}$. The D-brane assignments of the quark singlets U_a^c, D_a^c are more flexible in this model, since they can be of the form $C_i^{5_1}$, with $i = 1, \ldots, 3$. If the quark singlets of all three families are states of the type $C_3^{5_1}$ to obtain large top (and bottom) Yukawa couplings, we are led once again to the democratic texture in eq. (3.2). As shown in [26], this democratic texture is related to the more standard hierarchical texture with only the 33 entry being non-zero. The leading perturbations placed in the third row and column are related to the various hierarchical textures in the 23 block, and so on. Since this pattern of the quark mass matrices has been discussed at great length in [26] (to which we refer the reader for further details), we refrain from analyzing this scenario further. However, we do emphasize the rather interesting point that since the leading order Yukawa couplings (which are related to the gauge couplings at string tree-level) are equal, the permutation symmetry $S(3)_L \times S(3)_R$ emerges automatically from the theory.

However, since in the model above the quark singlets arise from a single D-brane sector, there is some flexibility in their possible D brane assignments. We have just seen how the democratic texture can arise if all quark singlets are $C_3^{5_1}$ fields; now we consider textures in the case in which only the quark singlets of the third family are states of this type. Certainly this is a minimal requirement in order to have a large top (and bottom) Yukawa coupling. Regarding the quark singlets of the first and second generations, there are several options; however, we prefer for now to investigate the case in which only the third family has large Yukawa couplings, such that $U_{1,2}^c, D_{1,2}^c \sim C_{1,2}^{5_1}$. In any case, with this assignment of the quark singlets (using eq. (2.2)), the MSSM quark superpotential couplings take the form

$$W_{\text{quark}} = \sum_{a,b=1}^{3} \delta_{b3}[Y_{ab}^u Q_a U_b^c H_u + Y_{ab}^d Q_a D_b^c H_d]\,, \tag{3.3}$$

and hence the (rescaled) Yukawa coupling matrices in flavor space of the effective lagrangian take the form of the *single right-handed democratic texture* referred to in eq. (1.1):

$$Y_{u,d} = \mathcal{O}(g_{5_1}) \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 1 \\ 0 & 0 & 1 \end{pmatrix}. \tag{3.4}$$

Since the up and down quark singlets are assumed to have identical D-brane assignments, the Yukawa matrices are identical to leading order. This in turn implies top-bottom unification at the string scale, and hence large $\tan\beta$ is required. This is sufficient for our purposes here; in a more general D-brane model one could envisage the two leading order Yukawa matrices to be controlled by two different gauge couplings of unequal strength, leading to lower values of $\tan\beta$. Note once again that an identical analysis can be carried out in the T-dual picture in which the SM gauge group is embedded within the nine-brane sector and one of the 5_i brane sectors. For example, one can consider the T-dual picture in which SU(3) and $\mathrm{U}(1)_Y$ originate from the nine-branes, and SU(2) from one of the 5_i sector. Identical Yukawa matrices to those in eq. (3.4) are obtained from the corresponding D-brane assignments: Q_a, L_a, $H_{u,d}$ are C^{95_1}, while U_3^c, D_3^c are C_1^9 fields.

Having been led to the form of the Yukawa matrix in eq. (3.4) it is necessary to address the question of whether it leads to a viable texture for quark masses and mixing angles. Surprisingly, even though the diagonalization of the u, d mass matrices separately involve large angle rotations, we find the CKM mixing angles are automatically small. The usual textures proposed for describing the quarks are based on hierarchical structures which are diagonalized by small mixing angles, and the form of the Yukawa matrices given above looks very radical compared to such schemes. In order to determine the desired pattern of perturbations in this case we use the analytic methods for diagonalizing a Yukawa matrix with a dominant coupling of a single right-handed fermion field developed in [29]. In diagonalizing the Yukawa matrices we may consider only the rotations of left-handed fields corresponding to unitary transformations on the left (the right-handed rotations always give negligible corrections to eigenvalues and left-handed rotation angles.) Each matrix may be approximately diagonalized by a product of left-handed rotations of the form $R^u_{12}R^u_{13}R^u_{23}Y^u$, $R^d_{12}R^d_{13}R^d_{23}Y^d$, corresponding to the angles $\theta_{12}, \theta_{13}, \theta_{23}$ familiar from the standard parameterization of the CKM matrix which involves a similar factorization. Then the CKM matrix is given by

$$V_{CKM} = {R^u_{12}}^\dagger {R^u_{13}}^\dagger {R^u_{23}}^\dagger R^d_{23} R^d_{13} R^d_{12} \,. \tag{3.5}$$

Let us write the Yukawa matrices in general as

$$Y_{u,d} = \begin{pmatrix} a'_{u,d} & a_{u,d} & d_{u,d} \\ b'_{u,d} & b_{u,d} & e_{u,d} \\ c'_{u,d} & c_{u,d} & f_{u,d} \end{pmatrix} , \tag{3.6}$$

where $a', b', c' \ll a, b, c \ll d, e, f$, then we find the mixing angles are given by [29]

$$\begin{aligned} \tan\theta^{u,d}_{23} &\approx \frac{e_{u,d}}{f_{u,d}} \qquad \tan\theta^{u,d}_{13} \approx \frac{d_{u,d}}{\sqrt{e^2_{u,d}+f^2_{u,d}}} \\ \tan\theta^{u,d}_{12} &\approx \frac{c^{u,d}_{13} a_{u,d} - s^{u,d}_{13}(s^{u,d}_{23} b_{u,d} + c^{u,d}_{23} c_{u,d})}{c^{u,d}_{23} b_{u,d} - s^{u,d}_{23} c_{u,d}} \,. \end{aligned} \tag{3.7}$$

Using these results and assuming the following Yukawa matrices

$$Y^u(0) = \mathcal{O}(g_{5_1}) \begin{pmatrix} \lambda^8 & \lambda^4 & 1+\mathcal{O}(\lambda) \\ \lambda^8 & \lambda^4 & 1+\mathcal{O}(\lambda^2) \\ \lambda^8 & \lambda^4 & 1+\mathcal{O}(\lambda^2) \end{pmatrix}, \; Y^d(0) = \mathcal{O}(g_{5_1}) \begin{pmatrix} \lambda^4 & \lambda^2 & 1+\mathcal{O}(\lambda) \\ \lambda^4 & \lambda^2 & 1+\mathcal{O}(\lambda^2) \\ \lambda^4 & \lambda^2 & 1+\mathcal{O}(\lambda^2) \end{pmatrix}, \tag{3.8}$$

where $\lambda \approx 0.22$ is the Wolfenstein parameter, we find an acceptable hierarchy of eigenvalues given by $m_u/m_t \sim \lambda^8$, $m_c/m_t \sim \lambda^4$, $m_d/m_b \sim \lambda^4$, $m_s/m_b \sim \lambda^2$, and a CKM matrix of the form

$$V_{CKM} \sim \begin{pmatrix} 1 & \lambda & \lambda^3 \\ \lambda & 1 & \lambda^2 \\ \lambda^3 & \lambda^2 & 1 \end{pmatrix}. \tag{3.9}$$

Note in eq. (3.8) that the small terms proportional to powers of λ involve unknown coefficients of order unity, whereas the leading order terms in the third column are equal. Therefore, from eqs. (3.7) and (3.8) it is clear for example that the 23 mixing angles in the u,d sectors are separately large but cancel to order λ^2 in the formation of the CKM matrix in eq. (3.5), since the 23 element of $R^{u}_{23}{}^{\dagger} R^d_{23}$ is

$$c^u_{23} s^d_{23} - s^u_{23} c^d_{23} = \frac{f_u e_d - e_e f_d}{\sqrt{e_u^2 + f_u^2}\sqrt{e_d^2 + f_d^2}} \sim \lambda^2 \tag{3.10}$$

due to the cancellation of the leading order terms. The remaining elements of the CKM matrix are obtained straightforwardly in a similar manner, to leading order in the perturbative expansion in λ, leading to the result in eq. (3.9). Note that the Yukawa matrices and CKM matrix given above are the values at the string scale, which is flexible within type-I string models. Assuming for simplicity that the string scale $M_{\rm String}$ and GUT scale $M_G \sim 3\times 10^{16}$ GeV coincide, renormalization group (RG) evolution from the string scale to the electroweak scale will affect the numerical values of these matrices; however, the form of the hierarchy is preserved by the RGE's. Note that if the D-brane assignments of the charged leptons and the right-handed neutrinos mirrored that of the quarks, the leading order coupling of the single right-handed fermion fields predicts single right-handed neutrino dominance which was introduced by one of us on phenomenological grounds in order to account for the concurrent large mixing angles and hierarchical neutrino masses [28]. However, the precise cancellation of the 23 mixing angles described above will also occur in the lepton sector. This is quite problematic for this model, so we defer discussion of the lepton sector until later.

Turning to the question of the structure of the smaller Yukawa couplings, several comments are in order. While the details of the D-brane assignments of the quark singlets of the first and second generations in this case does not affect the leading order structure of the Yukawa matrices, they will have an impact both on the possible mechanisms for generating the smaller Yukawa couplings and on the form of the soft breaking parameters. A possibility that has recently been explored [27] is that the index i labels the three different families. In this case the quark singlets take the form $U_1^c, D_1^c \sim C_1^{5_1}$, $U_2^c, D_2^c \sim C_2^{5_1}$, $U_3^c, D_3^c \sim C_3^{5_1}$. (One could imagine that the same assignments would hold for the leptons, such that $E_3^c \sim C_3^{5_1}$, $E_{1,2}^c \sim C_{1,2}^{5_1}$; the discussion of the lepton sector will be deferred to a later section). It is also important to note that the perturbations required in $Y_{u,d}$ above imply that the first and second generation quark singlets must be distinguished in some way, by D-brane assignments and/or U(1) charges. In addition, note that the 13 elements of both $Y_{u,d}$ must have corrections of $\mathcal{O}(\lambda)$ while the 23 and 33 elements must have corrections of $\mathcal{O}(\lambda^2)$, such that the true "democracy" of the couplings of the right-handed singlet of the third family must be broken by higher-order effects. At this stage we can not say much more about whether or not such perturbations are plausible, since the question of how to generate such small Yukawa couplings in the type-IIB orientifold models has not been fully addressed.

We close this section with a brief discussion of the form of the soft parameters at high energy in the D-brane model (2) with SU(3) and U(1)$_Y$ from the 5_1 sector and SU(2) from the 5_2 sector, assuming the dilaton and moduli are the primary mediators of supersymmetry breaking. This model was first studied in [19], based on the work of [18]. With the D-brane assignments of the MSSM fields as listed above, the form of the soft parameters are changed slightly; we assume further for simplicity that the D-brane assignments for the leptons are similar to those of the quarks, such that the Yukawa matrices have the single right-handed democratic structure studied above (and hence we are neglecting for now the problems associated with obtaining the large 23 mixing angle). To display the manifest symmetry of the low energy effective lagrangian as given in [18] between the various D-brane sectors, we utilize the following notation for the usual Goldstino angle parameterization [24, 25, 18]:

$$\begin{aligned} X_0 &= \sin\theta \\ X_1 &= \cos\theta\Theta_1 \\ X_2 &= \cos\theta\Theta_2 \\ X_3 &= \cos\theta\Theta_3\,, \end{aligned} \tag{3.11}$$

in which θ, $\Theta_{1,2,3}$ are the Goldstino angles which measure the relative contributions of the dilaton and moduli fields to SUSY breaking. In this notation $\sum_{i=0}^{3} X_i^2 = 1$.

In our model in which the SM gauge group is split between the 5_1 and 5_2 sectors,[5]

[5]For the T-dual picture in which the SM gauge group is split between the nine brane sector and one of the five-brane sectors, the soft breaking parameters follow from the expressions given below with the replacements $X_0 \to X_3$, $X_1 \to X_0$, $X_2 \to X_1$, $X_3 \to X_2$, as can be derived from the expressions for the effective action given in [18]. Hence, it is always possible to choose a set of VEV's of the dilaton and moduli fields in the T-dual picture which yield an identical spectrum (and hence identical phenomenology).

the gaugino masses $M_{1,2,3}$ are given by

$$\begin{aligned} M_1 &= \sqrt{3} m_{3/2} X_1 e^{-i\alpha_1} = M_3 \\ M_2 &= \sqrt{3} m_{3/2} X_2 e^{-i\alpha_2} \,, \end{aligned} \tag{3.12}$$

and the soft mass-squared parameters take the form

$$\begin{aligned} m^2_{Q_a} &= m^2_{L_a} = m^2_{H_{u,d}} = m^2_{3/2}\left(1 - \frac{3}{2}(X_0^2 + X_3^2)\right) \\ m^2_{U_1} &= m^2_{D_1} = m^2_{E_1} = m^2_{3/2}(1 - 3X_0^2) \\ m^2_{U_2} &= m^2_{D_2} = m^2_{E_2} = m^2_{3/2}(1 - 3X_3^2) \\ m^2_{U_3} &= m^2_{D_3} = m^2_{E_3} = m^2_{3/2}(1 - 3X_2^2) \,, \end{aligned} \tag{3.13}$$

in which $m_{3/2}$ is the gravitino mass parameter, and $\alpha_{1,2}$ are the (assumed) phases of the F-component VEV's of the moduli fields. For the trilinear couplings $\tilde{A}^{u,d,e}$ of the soft breaking lagrangian, in supergravity models $\tilde{A}^{u,d,e}_{ab}$ has the same hierarchical structure at the Yukawa matrix, but is not directly proportional to it, since each element differs from the corresponding element of the Yukawa matrix by a numerical coefficient of order unity. In our case, the general form of the trilinear couplings $\tilde{A}^{u,d,e}_{a,b}$ is given by

$$\tilde{A}^{u,d,e} \sim \begin{pmatrix} 0 & 0 & A_{u,d,e} Y_{13} \\ 0 & 0 & A_{u,d,e} Y_{23} \\ 0 & 0 & A_{u,d,e} Y_{33} \end{pmatrix} , \tag{3.14}$$

in which $-A_{u,d,e} = M_1 = M_3$. Regarding the generation of the remaining entries of $\tilde{A}_{u,d,e}$, there are two possibilities, depending on the mechanism utilized to generate the small Yukawa couplings. The first is that the effective Yukawa couplings can be generated at the string scale via nonrenormalizable operators (for example via the anomalous U(1)'s). The effective soft trilinear couplings can then be computed using the standard supergravity techniques; the soft trilinear couplings for this case have recently been presented in [27].[6] The other possibility is to generate the small Yukawa couplings at lower energy scales, in the effective quantum field theory; the soft trilinear couplings will also be generated if SUSY is broken at that scale. Due to our lack of knowledge as to how the small Yukawa couplings are generated, we will also not further speculate as to the form of the corresponding soft trilinear couplings.

Note that the spectrum above exhibits certain nonuniversalities in the values of the soft mass-squared parameters at the string scale, which must in general be checked to make sure that FCNC bounds are not violated. A detailed study of this

[6] Note that since [27] assumes arbitrary Yukawa matrices and yet utilize the D-brane structure to dictate the form of the soft trilinear couplings, [27] assumes implicitly that the structure of the trilinear couplings does not have to mirror the basic structure of the Yukawa couplings, contrary to what is expected within supergravity models.

question would include RG evolving the parameters to the electroweak scale, which is beyond the scope of this paper. However, it is amusing to note in this case that since the m_Q^2 parameters are universal and diagonal (since all quark doublets are $C^{5_1 5_2}$ states) at the string scale, the off-diagonal entries of the LL mass insertions at the electroweak scale will be well below the FCNC bounds given in [30]. Although the soft mass-squared parameters of the right-handed singlets are nonuniversal, the FCNC bounds on the RR mass insertions are much weaker, such that the nonuniversality among the right-handed fields is essentially unconstrained, as we have verified.

4. Lepton textures

Let us turn to a more systematic approach to the possible couplings in the lepton sector. In doing so, we note that there are several issues related to the necessary presence of right-handed neutrinos which complicate the analysis of the lepton sector considerably. First, it is clear that the potential presence of Majorana as well as Dirac mass terms for the neutrinos makes the interpretation of the phenomenological constraints on the lepton masses and mixings much more complicated. In addition, the embedding of the SM gauge group within the D-brane sectors does not provide much guidance in determining the D brane assignments of the right-handed neutrinos, since these fields are SM gauge singlets (in contrast, the charged leptons necessarily arise from open string sectors, in which one end of the open string must touch the brane sector which includes $\mathrm{U}(1)_Y$, while the other end can terminate on e.g. an anomalous U(1) from either the same set of branes or another brane sector). The possibility that the right-handed neutrinos are "bulk" fields (i.e., fields from the closed string sector) has been explored within more general "brane world" scenarios and can provide interesting alternatives to the traditional seesaw mechanism (see e.g. [31, 32] for recent explorations within particular string-motivated models, which however rely on the existence of at least one large extra dimension). For example, the right-handed neutrinos could be viewed as superpartners of moduli fields; the needed couplings of such fields to the left-handed neutrinos to generate Dirac mass terms could arise from nonperturbative operators. However, this approach is beyond the scope of this paper, and therefore we defer this question to a future study.

We thus restrict ourselves to the possibility that the right-handed neutrinos are open string states (see also [16]). However, it is important to note that the traditional view of right-handed neutrinos within GUT-type models, e.g. as part of the 16-plet in SO(10), or within $E(6)$ multiplets, is not compatible with the possible gauge groups and representations of the open string states within type-I models. In particular, it is possible to prove on general grounds that both exceptional groups and spinor representations of SO(10) cannot be generated through Chan-Paton charges [33]. However, the Pati-Salam subgroup $\mathrm{SU}(4) \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$ does arise in specific type-I constructions [15, 14] and we return to this later.[7] One can conclude that the origin of the right-handed neutrinos and their couplings within the type-IIB orientifolds is a complicated issue. Therefore, we do not attempt to exhaust all possibilities, but instead focus on illuminating possible constraints on the allowed couplings of these fields, within the assumption that they arise from the open string sectors of the theory.

We begin by stating the constraints on the mixings in the neutrino sector, as dictated by recent experimental results. Essentially, in order to have a large 23 angle (as required by SUPER-KAMIOKANDE) and a small 13 angle (as required by CHOOZ), we require the charged lepton Yukawa matrix Y_e and light neutrino Majorana matrix m_{LL} to take the leading order form

$$Y_e(0) = \mathcal{O}(1) \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \qquad m_{LL}(0) = \mathcal{O}(1) \begin{pmatrix} 0 & 0 & 0 \\ 0 & 1 & 1 \\ 0 & 1 & 1 \end{pmatrix}, \tag{4.1}$$

again assuming large $\tan\beta$ for simplicity. In the above recall that m_{LL} is the the matrix obtained after the seesaw mechanism,

$$m_{LL} = Y_n M_{RR}^{-1} Y_n^T \,, \tag{4.2}$$

in which M_{RR} is the heavy RH Majorana matrix, and the notation for the Dirac Yukawa couplings can be understood from the superpotential

$$W = \sum_{a,b} (Y_e)_{ab} L_a E_b^c H_d + \sum_{a,p} (Y_n)_{a,p} L_a N_p^c H_u + \sum_{p,q} (M_{RR})_{pq} N_p^c N_{q.}^c \,, \tag{4.3}$$

with M_{RR} given by the VEV of a singlet field Σ, such that $M_{RR} = Y_M \langle \Sigma \rangle$.

Before we begin, a comment about M_{RR} is in order. Due to the form of the allowed superpotential couplings in the type-IIB orientifold models (see e.g. eq. (2.2)), it appears to be quite difficult to obtain diagonal entries in the heavy Majorana mass matrix unless the RH neutrinos are all intersection states. This can be seen from the general form of the superpotential; for example, if N_3^c were a $C_1^{5_1}$ state, there is no term in the superpotential of the form $C_1^{5_1} C_1^{5_1} X$ (with X any other matter field, which presumably gets a large VEV to set the heavy Majorana mass scale). However, there are terms of the form $C^{5_1 5_2} C^{5_1 5_2} X$ for $X = C_3^{5_1}$ (and certainly $X \neq C^{5_1 5_2}$); hence diagonal Majorana terms (from trilinear superpotential couplings) appear to be possible only for right-handed neutrinos with $C^{5_1 5_2}$ assignments.

We first consider the possibilities for obtaining this form of Y_e and m_{LL} in the D-brane models given so far which lead to single right-handed democracy in the quark sector. Assuming the D-brane assignments for the charged leptons and the

[7] See also [7] for a recent study of (free fermionic) perturbative heterotic string models with the Pati-Salam gauge group.

right-handed neutrinos mirror that of the quark singlets such that $E_3^c, N_3^c \sim C_3^{5_1}$, the Yukawa couplings $Y_{e,n}$ take the form

$$Y_{e,n} = \mathcal{O}(g_{5_1}) \begin{pmatrix} 0 & 0 & 1 \\ 0 & 0 & 1 \\ 0 & 0 & 1 \end{pmatrix}, \tag{4.4}$$

which leads to single right-handed neutrino dominance (SRHND) [28]. However, although the potential viability of the SRHND mechanism has been demonstrated in [28, 29], in this case there are two difficulties, such that these couplings do not lead to the desired form of Y_e and m_{LL}. First, the presence of the large 13 elements is problematic; this element needs to vanish to leading order for the physical 13 mixing angle to be small. Second, in [28] m_{LL} and Y^e may both have large 23,33 elements leading to a large physical 23 mixing angle since the Yukawa elements are only *approximately* of order unity. This is in sharp contrast to the D-brane model here, in which the precise equality between the 23 and 33 entries of the Yukawa matrices (at least at string tree-level) leads to a natural cancellation of the physical 23 mixing angle, as was shown above for the quark sector. In general, what one finds for m_{LL} in this case is the following:

$$m_{LL} \simeq (M_{RR}^{-1})_{33} \begin{pmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{pmatrix}, \tag{4.5}$$

which is clearly incompatible with phenomenology. Furthermore, the fact that the seesaw mechanism selects out the 33 element of M_{RR}^{-1} may also be problematic, since this term is not generated at the trilinear order in this model (as explained above).

Although the main purpose of this paper is to explore flavor physics in the type-I framework without resorting to family symmetries, one could imagine utilizing family symmetries (e.g. from the anomalous U(1)'s generically present in this class of superstring models) to obtain the desired form of Y_e and m_{LL}.[8] We conclude that it appears to be quite difficult to achieve viable neutrino textures in this model in the absence of additional suppression mechanisms because of the way that the SM gauge group embedding treats all doublets symmetrically (although this led to interesting

[8]In the case of a single U(1) family symmetry for simplicity, the suppression of the 13 and 23 entries of Y_e can be obtained by requiring $Q_{L_1,L_2} \neq Q_{L_3} = -(Q_{E_3} + Q_{H_d})$. For example [29], a simple possibility is to assume charges of the form $Q_{L_1} = -1$, $Q_{L_2} = 1$, $Q_{L_3} = 0$, $Q_{N_{1,2}} = 1/2$, $Q_{N_3} = -1/2$, $Q_{E_1} = 5$, $Q_{E_2} = 1$, $Q_{E_3} = 0$. In this case $(Y_e)_{33}$ is allowed but $(Y_e)_{13}$ and $(Y_e)_{23}$ are forbidden at leading order. Turning to the neutrinos, $(Y_n)_{13} = Q_{L_1} + Q_{N_3} = -3/2$, $(Y_n)_{23} = Q_{L_2} + Q_{N_3} = 1/2$, and $(Y_n)_{33} = Q_{L_3} + Q_{N_3} = -1/2$; we see that the 13 element of Y_n is suppressed by one unit of charge (i.e. power of λ) relative to the 23 and 33 which are of the same order. Since the half integer power of λ may be factored out, this leads to the desired form for Y_n (overall factors do not affect the mixing angles). A complete analysis is beyond the scope of this paper, and we do not explore such possibilities further.

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features in the quark sector), irrespective of the form of the heavy Majorana mass matrix M_{RR} (which may be difficult in general to generate, at least for the case in which this term arises from a trilinear coupling). Note that this is an example of how restrictive such string theories can be.

One may ask whether or not there exist alternative D-brane embeddings of the SM in which this type of lepton texture can emerge gracefully from the underlying structure of the theory, rather than relying on the U(1) symmetry. The main point to keep in mind regarding the lepton sector is that the lepton doublets $L_{2,3}$ of the second and third families must be distinguishable from each other from the point of view of the D brane assignments in order to have leading-order Yukawa couplings of the desired form. With this structure in mind, we will present a representative model of this type which can emerge from the splitting of the SM gauge groups among at most two D-brane sectors for simplicity. While there may be greater flexibility in considering models in which the SM gauge group consists of linear combinations of gauge groups from more than two D-brane sectors, such an analysis is beyond the scope of this paper. Therefore, we consider the general case of type-IIB orientifold models in which the SM gauge group is split between two D-brane sectors, and allow both SU(2) and $\mathrm{U}(1)_Y$ to be linear combinations of gauge groups from both of the D-brane sectors, in the spirit of [15]. In this case, by necessity we let the phenomenological constraints on allowed couplings in the lepton sector (rather than the SM gauge group embedding) dictate the D-brane assignments of the MSSM fields.

Considering the case in which the SM gauge group is split between two five-brane sectors, one essentially needs to assume that $\mathrm{SU}(2)_L$ and $\mathrm{U}(1)_Y$ are linear combinations of gauge groups from both branes such that the lepton doublets may have the D-brane assignments $L_2 = C_3^{5_2}$, $L_3 = C_3^{5_1}$. To obtain the desired Yukawa structure, one also needs to distinguish between the MSSM Higgs doublets $H_{u,d}$, for example as $H_u = C^{5_1 5_2}$, $H_d = C_1^{5_1}$. Then, if the right-handed neutrinos and the charged lepton singlets of the third family are fields of the type $N_3 = C^{5_1 5_2}$, $E_3 = C_2^{5_1}$ (and $N_{1,2}, E_{1,2} \neq C^{5_1 5_2}, E_3 = C_2^{5_1}$), the Yukawa couplings matrices in the lepton sector take the form

$$Y_e = \mathcal{O}(g_{5_1}) \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & 1 \end{pmatrix}, \qquad Y_n = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & \mathcal{O}(g_{5_2}) \\ 0 & 0 & \mathcal{O}(g_{5_1}) \end{pmatrix}, \tag{4.6}$$

The heavy Majorana matrix can then take the form to leading order (assuming that there is a singlet field $\Sigma = C_3^{5_1}$ (or $C_3^{5_2}$)),

$$M_{RR} = \begin{pmatrix} 0 & 0 & 0 \\ 0 & 0 & 0 \\ 0 & 0 & O(g_{5_{1,2}}) \equiv 1/c \end{pmatrix}, \tag{4.7}$$

which leads to

$$m_{LL} \sim \begin{pmatrix} 0 & 0 & 0 \\ 0 & cg_{5_2}^2 & cg_{5_1}g_{5_2} \\ 0 & cg_{5_1}g_{5_2} & cg_{5_1}^2 \end{pmatrix}, \tag{4.8}$$

which is similar to the desired form. Note that this is in fact quite nontrivial; the D-brane assignment $N_3^c = C^{5_1 5_2}$ was crucial in obtaining the desired diagonal entry in the heavy Majorana mass matrix. Since gauge coupling unification is not automatic in these models, these entries are not in general guaranteed to be equal at the string scale (in any event RG effects will likely spoil any exact equality of these entries at the string scale upon evolution to the electroweak scale). It does not appear to be possible to construct a model within this framework with this asymmetry between the 23 elements of the neutrino and charged lepton sectors and exact degeneracy between the 23 and 33 elements of the neutrino Yukawa matrix.[9] Also note that eq. (4.6) maintains the simple principle of SRHND in that a single right-handed neutrino couples at leading order, leading to a mass matrix in eq. (4.8) with vanishing 23 sub-determinant which enforces a natural 23 mass hierarchy [28].

For completeness, let us consider now the possible quark Yukawas in this framework. With the D-brane assignment of the Higgs doublets $H_{u,d}$ as above, it is necessary that (at least for the third family in each case) the quark doublets are of the form $C_3^{5_1}$, up-type quark singlets are intersection states $C^{5_1 5_2}$, and the down-type quark singlets are states of the type $C_2^{5_1}$ in order to have nonzero Yukawa matrices at leading order. In this case, the right-handed democracy only emerges if one assumes that the quark doublets of all three families are $C_3^{5_1}$ states (which may or may not be plausible), and does not follow intrinsically from the underlying structure of the model. It is also clear from the quark assignments that the most natural interpretation is that SU(3), along with SU(2) and U(1)$_Y$, is some linear combination of gauge groups from both sectors. (One could imagine the slightly more artificial scenario in which SU(3) arises only from the 5_1 sector. In this case the up-type quark singlets U_i would have to be singlets under the gauge group of the 5_2 sector).

While we present this model as an existence proof that D-brane assignments can in principle allow for lepton Yukawa matrices of the form in eq. (4.1), it is important to note that this model is somewhat complicated, and it is not clear at all whether or not this type of model could arise within explicit string models. We must therefore emphasize that such structures in which the lepton doublets are treated asymmetrically in general do not appear to emerge in a graceful manner from the point of view of the underlying type-I theory, but at least they can be accommodated in this framework.

Finally we discuss the prospects for understanding the origin of right-handed neutrino masses in D-brane models which do not directly lead to the standard model

[9] Other possible models which allow the general structure of eq. (4.1) are basically related to the form of this model by trivial permutations of possible D-brane assignments for the MSSM fields.

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gauge group, but instead to some larger group which includes $\mathrm{SU}(2)_R$, such as the Pati-Salam models originating from the type-I constructions in [15], [14]. The Pati-Salam gauge group was also considered some time ago in the framework of four-dimensional fermionic string constructions [34]. Under the Pati-Salam gauge group $\mathrm{SU}(4) \times \mathrm{SU}(2)_L \times \mathrm{SU}(2)_R$ a complete quark and lepton family is accommodated in the reducible representation $F(4,2,1) + F^c(\bar{4},1,\bar{2})$, and the gauge group is broken by heavy Higgs $H(4,1,2)$ which develop VEV's in the neutral component. Right-handed neutrino masses originate from the couplings to gauge singlets $\phi(1,1,1)$ and $\theta(1,1,1)$: $F^cH\phi$ and $\phi\phi\theta$. When θ gets a VEV this leads to a ϕ mass term. Below the ϕ mass the renormalizable couplings induce the nonrenormalizable operator F^cHHF^c. Finally when H gets a VEV this operator induces a mass term for the right-handed neutrinos contained in F^c. The phenomenology of neutrino masses based on this mechanism was discussed in the early references [35], and post-Super-Kamiokande work is presently in progress [36]. In order to implement this mechanism in the D-brane framework, many of the same issues arise for the coupling of the singlet fields θ, ϕ as for the previous discussion in which it was assumed that the right-handed neutrino was a gauge singlet.

5. Conclusions

In this paper, we have presented a first analysis of possible (leading order) Yukawa textures within four-dimensional type-I superstring-motivated models in which the SM gauge group is split between different D-brane sectors. Due to the important role of the intersecting D branes, such models can have phenomenological implications which are quite distinctive compared with that of traditional perturbative heterotic superstring models, or models in which all three SM gauge groups are embedded within a single D-brane sector. For example, such models have inspired previous interest in the context of CP violation due to the tree-level nonuniversality of the gaugino masses dictated by the SM gauge group embedding. Here we have demonstrated that the leading order Yukawa matrices for the quarks and leptons can have novel structures which are also dictated by the way in which the SM gauge group is split between the different D-brane sectors, such that the D-brane assignments of the matter fields can either replace or supplement family symmetries. In certain models, we found that the leading order structure of the Yukawa matrices has a novel form which we label "single right-handed democracy", which can be a viable texture for the quark sector. The origin of the right-handed neutrinos poses significant challenges within this framework, and is an interesting avenue for future study.

Although it has been possible to study many interesting issues regarding the flavor structure of the Yukawa couplings, we emphasize that an analysis such as the one presented in this paper is simply the beginning stage in the investigation of flavor physics within the type-I string picture. The first step is of course to continue

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to construct quasi-realistic string-derived models to explore the possible ways in which the SM can be embedded within string theory. It is encouraging to note that there has been recent progress along these lines. For the purposes of this particular study, it is crucially important to determine if the structures outlined in this paper can arise within explicit string models. However, much remains to be done even within the type-I string-motivated approach. Since only the trilinear superpotential couplings at the string scale (which have Yukawa couplings of $\mathcal{O}(1)$) are available, it is only possible to determine the leading order structure of the Yukawa matrices for the quarks and leptons, and it is not possible yet to address the issue of how the smaller Yukawa couplings are generated from higher-dimensional operators. Hence, the important questions of whether or not the perturbations to the leading-order Yukawas lead to viable fermion mass hierarchies and whether or not a sizeable CKM phase is generated are beyond the scope of this paper, and certainly worthy of future exploration.

To conclude, we believe that the type-I string framework provides a rich and promising setting for exploring novel phenomenological consequences of superstring theory. This analysis in particular demonstrates the ways in which this framework can provide new ways to think about flavor physics, which may help illuminate the elusive resolution to the flavor puzzle.

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