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Discrete symmetries in semi-realistic orientifold compactifications

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This proceeding is based on arXiv:1211.1017 [hep-th] where we study the presence of discrete gauge symmetries in D-brane semi-realistic quivers. After establishing the constraints on the transformation behaviour of the chiral matter for the presence of a discrete gauge symmetry we perform a systematic search for discrete gauge symmetries within semi-realistic local D-brane realizations, based on four D-brane stacks. This search reveals that Proton hexality, a discrete symmetry which ensures the absence of R-parity violating terms as well as the absence of dangerous dimension 5 proton decay operators, is only rarely realized. Furthermore, we do not find any semi-realistic local D-brane configurations exhibiting a family dependent discrete gauge symmetry.

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1. Introduction

The Minimal Supersymmetric Standard Model (MSSM) provides solutions to some open issues of the Standard Model (SM), such as the hierarchy problem, exhibiting a natural dark matter candidate (i.e. the lightest supersymmetric particle, LSP) as well as gauge coupling unification. On the other hand, it exhibits some severe phenomenological problems, allowing renormalizable R-parity breaking operators consistent with supersymmetry and gauge invariance of the superpotential that do lead to a disastrous effects like high proton decay rate.

More precisely, the renormalizable MSSM gauge invariant superpotential terms are

$$W_{MSSM} = Y_U\, Q_L U_R H_u + Y_D\, Q_L D_R H_d + Y_L\, L E_R H_d + \mu\, H_u H_d \\ + \lambda_1\, U_R D_R D_R + \lambda_2\, Q_L L D_R + \lambda_3\, L L E_R + \alpha\, L H_u \,, \tag{1.1}$$

where the terms in the first line are the Yukawa couplings giving mass to quarks and leptons after electroweak symmetry breaking as well as the μ-term. The second line contains terms, so called R-parity violating terms, that do not conserve baryon and lepton number. They can lead to rapid proton decay, rendering the LSP unstable and thus eliminating the possibility of any SUSY particle being the dark matter candidate. Moreover, SM gauge invariance allows also for the dimension 5 proton decay operators

$$Q_L Q_L Q_L L \qquad\qquad U_R U_R D_R E_R \,, \tag{1.2}$$

which if not suppressed lead to a disastrous high proton decay rate.

In order to avoid these phenomenologically dangerous terms, there should be a mechanism within the MSSM that removes them. This role is usually played by discrete symmetries. For example, R-parity or proton triality B_3 remove all or some of the terms in the second line of (1.1). On the other hand proton hexality P_6, a discrete $\mathbf{Z}_6$ symmetry, forbids the R-violating terms and the dangerous dimension 5 operators.

Even if such symmetries solve these problems, their fundamental origin is unknown. Since *global* discrete symmetries are expected to be violated in consistent theories that includes quantum gravity [1–5]. An exception to that rule are discrete symmetries that have gauge origin aka *discrete gauge symmetries*. For instance abelian discrete symmetries $\mathbf{Z}_N$ are remnants of continuous $U(1)$ symmetries that are broken by scalars with charge N under the respective $U(1)$ acquiring vev's. However, the presence of a discrete symmetry seems fine-tuned unless there is a dynamical reason for the scalar field with charge N to acquire an appropriate vev. This possibility has been explored in several works in the past [6–9] and embedded in the context of heterotic string model building [10, 11].

Recently, in [12, 13] the authors explore the possibility of an even more fundamental/stringy origin of these discrete gauge symmetries as remnants of abelian gauge symmetries living on D-branes. Here we present the results of work [14] where we study the presence of discrete gauge symmetries in a class of promising MSSM D-brane quivers. After establishing the constraints on the transformation behaviour of the chiral matter for the presence of discrete gauge symmetries we perform a systematic search for discrete gauge symmetries within a class of local semi-realistic D-brane configurations.

	Q_L	U_R	D_R	L	E_R	N_R	H_u	H_d
A	0	0	-1	-1	0	1	0	1
L	0	0	0	-1	1	1	0	0
R	0	-1	1	0	1	-1	1	-1
$Q_{discrete}$	0	$-m$	$m-n$	$-n-p$	$m+p$	$n+p-m$	m	$-m+n$

Table 1: The family independent generators and the charges of the discrete $\mathbf{Z}_N$ gauge symmetries in the MSSM.

This proceeding is organised as follows: In section 2 we review the discrete gauge symmetries in the MSSM, using four-dimensional discrete anomaly conditions. In section 3 we discuss the constraints on the transformation behaviour of chiral matter that arise from string consistency conditions. Moreover, we establish the conditions on the transformation behaviour of the matter fields for the presence of a discrete gauge symmetry in D-brane compactifications. In section 4 we describe the systematic bottom-up search that we have perform by imposing the constraints for the presence of a discrete gauge symmetry, studied before, for a class of intriguing local D-brane configurations that exhibit a (semi-) realistic phenomenology. We present a characteristic example and the discrete symmetries that can appear as well as their phenomenological implications.

2. Discrete gauge symmetries in the MSSM from a field theory perspective

In this section, we review all family independent (non-R) discrete gauge symmetries of MSSM [6]. that satisfy the four-dimensional discrete gauge anomaly constraints, i.e. the mixed and gravitational anomalies

$$\mathscr{A}_{SU(3)SU(3)\mathbf{Z}_N} \quad , \qquad \mathscr{A}_{SU(2)SU(2)\mathbf{Z}_N} \quad , \qquad \mathscr{A}_{GG\mathbf{Z}_N} \; . \tag{2.1}$$

as well as allow for the Yukawa couplings

$$Q_L H_d D_R \qquad\qquad Q_L H_u U_R \qquad\qquad L H_d E_R \; . \tag{2.2}$$

Already shown in [7, 8] any family independent discrete gauge symmetry $\mathbf{Z}_N$ of the MSSM with generator g_N can be expressed in terms of products of powers of three mutually commuting generators A_N, L_N and R_N, i.e.

$$g_N = A_N^n \times L_N^p \times R_N^m \,, \qquad\qquad m,n,p = 0,1,...N-1 \; . \tag{2.3}$$

In table 1 we provide the charges of all chiral matter of MSSM under these three generators.

Solving (2.1, 2.2) one finds a finite class of solutions that contains of only $\mathbf{Z}_2$, $\mathbf{Z}_3$, $\mathbf{Z}_6$, $\mathbf{Z}_9$ and $\mathbf{Z}_{18}$ symmetries [6]. We present the results of all possible family independent discrete gauge symmetries of the MSSM in table 2. A similar analysis can be carried out for the MSSM with three additional right-handed neutrinos accompanied with a Dirac mass term. In that case one finds that all $\mathbf{Z}_2$, $\mathbf{Z}_3$, $\mathbf{Z}_6$ discrete gauge symmetries can be realized, however beyond those no further discrete gauge symmetry that can satisfy the discrete gauge anomaly conditions.

For a given discrete gauge symmetry, and for a specific choice of the parameters m, n and p we can easily evaluate the discrete charges for each MSSM field and concequently which terms are

N	n	p	m	Discrete gauge symmetries
2	0	0	1	R_2
3	0	0	1	R_3
	0	1	$(0,1,2)$	L_3, L_3R_3, $L_3R_3^2$
6	0	0	1	R_6
	0	2	$(1,3,5)$	$L_6^2R_6$, $L_6^2R_6^3$, $L_6^2R_6^5$
9	3	1	$(2,5,8)$	$A_9^3L_9R_9^2$, $A_9^3L_9R_9^5$, $A_9^3L_9R_9^8$
18	6	2	$(1,7,13)$	$A_{18}^6L_{18}^2R_{18}$, $A_{18}^6L_{18}^2R_{18}^7$, $A_{18}^6L_{18}^2R_{18}^{13}$

Table 2: All fundamental discrete gauge symmetries in the MSSM satisfying the anomaly cancellation conditions [6]. Here one allows for heavy fermions with fractional charges.

coupling	R_2	L_3R_3	R_3	L_3	$L_3R_3^2$	$L_6^2R_6^5$	R_6	$L_6^2R_6^3$	$L_6^2R_6$	$\mathbf{Z}_9$ & $\mathbf{Z}_{18}$
H_uH_d	✓	✓	✓	✓	✓	✓	✓	✓	✓	
LH_u		✓								
LLE_R		✓								
Q_LLD_R		✓								
$U_RD_RD_R$				✓						
$Q_LQ_LQ_LL$	✓		✓				✓			
$U_RU_RD_RE_R$	✓		✓				✓			
LH_uLH_u	✓	✓				✓				
N_RN_R	✓	✓				✓				

Table 3: Allowed superpotential terms for the respective discrete gauge symmetries [6].

excluded from the superpotential 1.1. The results are given in table 3. The $\mathbf{Z}_2$ symmetry R_2 is the usual matter parity [15] while L_3R_3 is Baryon triality [7]. Proton hexality, basically the product of matter parity and Baryon triality, is given by $L_6^2R_6^5$ and forbids all R-parity violating terms as well as the dangerous dimension 5 proton decay operators while still allowing for a μ-term H_uH_d and the Weinberg operator LH_uLH_u.

The above discussion on the allowed couplings for the respective discrete gauge symmetry applies specifically to the MSSM. Allowing for additional singlets, such as right-handed neutrinos, which do not acquire any vev does not change the analysis. However, the presence of right-handed neutrinos accompanied with a Dirac neutrino mass term raises the issue of the generation of small neutrino masses. A particular intriguing mechanism is the see-saw mechanism that requires large Majorana mass terms for the right-handed neutrinos. In the last line of table 3 we display which of the discrete symmetries permits for a Majorana mass term and thus allows the generation of small neutrino masses via the see-saw mechanism.

Finally, there exist two additional classes of discrete gauge symmetries, namely non-abelian discrete gauge symmetries and discrete R-symmetries. As recently pointed out the latter may play a special role in GUT theories, realised as a $\mathbf{Z}_4^R$ symmetry that forbids all R-parity violating terms as well as dimension 5 proton decay operators [16–18]. On the other hand non-abelian discrete gauge symmetries are often times invoked explaining various observations in flavour physics (see

Representation	Multiplicity
$\square\!\square_a$	$\#(\square\!\square_a) = \frac{1}{2}\left(\pi_a \circ \pi'_a - \pi_a \circ \pi_{O6}\right)$
$\begin{smallmatrix}\square\\\square\end{smallmatrix}_a$	$\#(\begin{smallmatrix}\square\\\square\end{smallmatrix}_a) = \frac{1}{2}\left(\pi_a \circ \pi'_a + \pi_a \circ \pi_{O6}\right)$
$(\square_a, \overline{\square}_b)$	$\#(\square_a, \overline{\square}_b) = \pi_a \circ \pi_b$
$(\square_a, \square_b)$	$\#(\square_a, \square_b) = \pi_a \circ \pi'_b$

Table 4: Chiral spectrum of intersection D-branes.

e.g. [19]). In this work we perform a systematic bottom-up D-brane analysis which ignores any specifics of the internal geometry. However non-abelian discrete gauge symmetries as well as discrete R-symmetries do rely on the details of the compactification manifold. Thus here we focus only on the subset of abelian discrete gauge symmetries.

3. Discrete symmetries in D-brane compactifications

D-brane model building provides an intriguing framework for semi-realistic model building. In those constructions the Standard Model lives on some stacks of branes[1]. Gauge fields are strings with both ends on the same stack and the chiral matter lives at the intersections. Here, we will focus on Type IIA constructions with intersecting D6 branes which wrap three-cycles π_x in the internal manifold. Each stack gives rise to an $U(N) = SU(N) \times U(1)$ gauge theory. All irreducible anomalies are cancelled via the *tadpole cancellation*, which ensures consistency and stability of the configuration:

$$\sum_x N_x \left(\pi_x + \pi'_x\right) = 4\pi_{O6} \ . \tag{3.1}$$

The sum runs over all D-brane stacks x in the given global setup and π'_x denotes the orientifold image cycle of π_x while π_{O6} denotes the orientifold cycle. We introduce a basis of three-cycles $\{\alpha_k\}$ and $\{\beta_k\}$ that are even and odd under the orientifold action, respectively, with $k = 1, ..., h_{21} + 1$. The choice of basis is such that $\alpha_k \cdot \beta_l = \delta_{kl}$ and $\alpha_k \cdot \alpha_l = \beta_k \cdot \beta_l = 0$. Then a three-cycle π_x and its orientifold image π'_x wrapped by a D-brane stack and its image D-brane stack, respectively, can be expanded in terms of this basis

$$\pi_x = \sum_k (m_x^k \alpha_k + n_x^k \beta_k) \qquad\qquad \pi'_x = \sum_k (m_x^k \alpha_k - n_x^k \beta_k) \ , \tag{3.2}$$

where m_x^k and n_x^k are integer and are usually referred to as wrapping numbers.

Multiplying 3.1 with the three-cycle $\pi_a$that is wrapped by the D-brane stack a, using table 4 one obtains after a few some manipulations [38–40]

$$\sum_{x \neq a} N_x \Big(\#(\square_a, \overline{\square}_x) + \#(\square_a, \square_x)\Big) + (N_a - 4)\#(\begin{smallmatrix}\square\\\square\end{smallmatrix}_a) + (N_a + 4)\#(\square\!\square_a) = 0 \ , \tag{3.3}$$

[1]For recent reviews on D-brane model building, see [20–23]. The first local bottom-up constructions were discussed in [24–26]. For original work on globally consistent non-supersymmetric intersecting D-branes, see [27–30], and for chiral globally consistent supersymmetric ones, see [31, 32]. For supersymmetric MSSM realizations, see [33–35], and for supersymmetric constructions within type II RCFT's, see [36, 37].

which is a constraint for each D-brane stack a of the D-brane setup. Due to the absence of antisymmetric representations for abelian gauge symmetries for a $U(1)$ stack, for a single D-brane stack, the constraint takes the form

$$\sum_{x\neq a} N_x\Big(\#(\square_a,\overline{\square}_x)+\#(\square_a,\square_x)\Big)+5\#(\square\!\square_a)=0 \qquad \mathrm{mod}\, 3\,. \tag{3.4}$$

Usually, the abelian part of the groups living on the branes are anomalous and gets a mass via the Stückelberg mechanism cancelling all the mixed anomalies [41–49]. However for the realization of the MSSM, it is required that at least one linear combination

$$U(1)=\sum_x q_x U(1)_x\,, \tag{3.5}$$

of these abelian factors originating from each stack of branes remains anomaly free (and massless) and does play the role of the hypercharge. The *masslessness condition* reads [25]:

$$\frac{1}{2}\sum_x q_x N_x(\pi_x-\pi_x')=0\,. \tag{3.6}$$

which after analogous manipulations as performed above turns into a constraint on the transformation behaviour of the chiral matter fields that reads

$$\frac{1}{2}\sum_{x\neq a} q_x N_x\#(\square_a,\overline{\square}_x)-\frac{1}{2}\sum_{x\neq a} q_x N_x\#(\square_a,\square_x)=\frac{q_a N_a}{2(4-N_a)}\left(\sum_{x\neq a} N_x\Big(\#(\square_a,\overline{\square}_x)+\#(\square_a,\square_x)\Big)+8\#(\square\!\square_a)\right) \tag{3.7}$$

where we have substituted the antisymmetrics that appear by using the tadpole condition 3.3.

The equations (3.1) and (3.6) are conditions on the three-cycles the D6-branes wrap, and imply the transformation behaviour of the four-dimensional chiral matter under the D-brane gauge symmetries. More specifically, the chiral matter fields cannot be distributed arbitrarily at the intersections of stacks of D-branes, but they have to obey the above conditions. Let us mention that the constraints (3.3), (3.4) and (3.7) are only necessary constraints but not sufficient. So any global D-brane construction has to satisfy those constraints, however a local D-brane configuration satisfying (3.3), (3.4) and (3.7) may not have a global realization.

3.1 Discrete Gauge Symmetries

For a discrete gauge symmetry $\mathbf{Z}_N$ arising from a linear combination

$$\mathbf{Z}_N=\sum_x k_x U(1)_x \tag{3.8}$$

to survive in the low energy effective field theory it has to satisfy:

$$\frac{1}{2}\sum_x k_x N_x(\pi_x-\pi_x')=0 \quad \mathrm{mod}\, N\,. \tag{3.9}$$

which represents a separate constraint for each D-brane stack a. Note that due to the non-integer prefactor $\frac{1}{2}$ in equation 3.9 the k_x do lie in the interval $(0,2N-1)$. After few manipulations and by using the tadpole condition (3.3) to eliminate the antisymmetrics, it becomes:

$$\frac{1}{2}\sum_{x\neq a} k_x N_x \#(\Box_a,\overline{\Box}_x) - \frac{1}{2}\sum_{x\neq a} k_x N_x \#(\Box_a,\Box_x) - \frac{k_a N_a}{2(4-N_a)}\left(\sum_{x\neq a} N_x\Big(\#(\Box_a,\overline{\Box}_x)+\#(\Box_a,\Box_x)\Big)+8\#(\Box\!\Box_a)\right)=0 \mod N . \tag{3.10}$$

One has to be slightly careful in using the tadpole constraint (3.1) to replace the antisymmetrics due to the fact that generically the prefactor is non-integer, in particular, in this case where the left hand side is not 0 but rather $0 \mod N$. One can compensate that by enlarging the interval for the k_x or by requiring an additional constraint arising from multiplying the homology class of the orientifold plane with the discrete symmetry constraint (3.8) [2]. This additional constraint reads

$$\sum_a k_a N_a \left(\#\Big(\begin{smallmatrix}\Box\\ \Box\end{smallmatrix}_a\Big) - \#(\Box\!\Box_a)\right) = 0 \mod N , \tag{3.11}$$

which after replacing the antisymmetrics in order not to have to distinguish between non-abelian and abelian D-brane stacks takes the form

$$\sum_a \frac{k_a N_a}{4-N_a}\left(\sum_{x\neq a} N_x\Big(\#(\Box_a,\overline{\Box}_x)+\#(\Box_a,\Box_x)\Big)+2N_a\#(\Box\!\Box_a)\right)=0 \mod N . \tag{3.12}$$

Let us mention that the constraints (3.10) and (3.12) do imply the vanishing of the various discrete gauge anomalies, such as $\mathscr{A}_{SU(N)SU(N)\mathbf{Z}_N}$ or $\mathscr{A}_{GG\mathbf{Z}_N}$. However, analogously to the abelian gauge symmetry these string theory constraints are more severe than just four-dimensional discrete gauge anomaly cancellation. Finally, again the constraints (3.10) and (3.12) do provide only necessary conditions, but not sufficient ones.

4. Systematic bottom-up search

In the work [14], we performed a systematic bottom-up D-brane analysis among promising semi-realistic local D-brane configurations found in [39, 50, 51][3] in which the Standard Modelis realized on 4 stacks of D-branes. Those promising quivers satisfy the consistency conditions (3.3), (3.4) and (3.7) as well as several phenomenological criteria. Among those rank that the desired Yukawa couplings giving masses to the fermions are realized perturbatively or non-perturbatively (where D-instantons play a very important role [68–71]) and R-parity violating terms as well as dangerous dimension 5 proton decay operators do not share the same quantum numbers as the desired Yukawa couplings.

[2]Note that for the abelian gauge symmetry such an additional constraint is not necessary, since one can use the tadpole constraint to replace the homology class of the orientfold plane by all the three-cycles wrapped by the D-brane stacks.

[3]For recent analogous work on semi-realistic bottom-up searches, see [40, 52–67].

Here we will study those local D-brane configurations with respect to discrete symmetries. We will analyse what quivers do satisfy the constraints to exhibit discrete symmetries and investigate their implications on the superpotential couplings.

Let us lay out the details of the search: For a chosen N we check whether a given linear combination of $U(1)$'s in terms of the vector (k_a,k_b,k_c,k_d), with the k_x's being integers, does satisfy the constraints (3.10) and (3.12). Due to the prefactor $\frac{1}{2}$ in eq. (3.10) we let the k_x run from 0 to $2N-1$.

Via a hypercharge shift we can find to any given solution (k_a,k_b,k_c,k_d) an additional equivalent solutions by adding the hypercharge. Thus $(k_a+my_a,k_b+my_b,k_c+my_c,k_d+my_d)$ is also a solution to the constraints (3.10) and (3.12) where m is an integer and the y_x denote the integer hypercharge embedding coefficients. In order to avoid overcounting we fix the discrete charge of Q_L for one family to be 0 by choosing $k_a = k_b$[4]. Thus we run only over three free integer parameter, namely k_a, k_c and k_d.

Additionally, we demand that the discrete symmetries allow for the quark and lepton Yukawa couplings in the superpotential, whose presence is crucial for the generation of low energy fermion masses. It turns out that this requirement is very stringent and rules out various discrete symmetries which otherwise satisfy the discrete top-down constraints (3.10) and (3.12).

Finally, we often find solutions for discrete gauge symmetries of higher degree due to the $\frac{1}{2}$ in (3.10) and (3.12), such as $\mathbf{Z}_{12}$, which eventually after determining the matter field charges turn out to be of lower degree from a pure MSSM point of view, since all matter charges have a common divisor. We take this into account when identifying the discrete symmetries but nevertheless display the linear combinations describing the discrete gauge symmetries in the D-brane language. Therefore, it frequently happens that $\mathbf{Z}_6$ symmetries contain coefficients that are higher than 12.

In [14] we investigated all of the promising four-stack quivers found in a systematic bottom-up search performed in [39, 50, 51]. Here we will present only the analysis for a specific hypercharge embedding. We will encounter two examples where the allowed discrete gauge symmetries forbid R-parity violating terms as well as the dangerous dimension 5 proton decay operators.

4.1 An example: Hypercharge $U(1)_Y = -\frac{1}{3}U(1)_a - \frac{1}{2}U(1)_b + U(1)_d$

We focus in a class of models where the hyper charge embedding is given by the $U(1)_Y = -\frac{1}{3}U(1)_a - \frac{1}{2}U(1)_b + U(1)_d$. In these models, we have the the MSSM particles plus three neutrinos where the spectra are displayed in table 5.

#	Q_L	D_R	U_R		L	E_R		N_R		H_u	H_d	
	$(\square_a,\overline{\square}_b)$	$(\overline{\square}_a,\overline{\square}_c)$	$(\overline{\square}_a,\overline{\square}_d)$	$\begin{smallmatrix}\square\\ \square\end{smallmatrix}_a$	$(\square_b,\overline{\square}_c)$	$(\square_c,\square_d)$	$\overline{\begin{smallmatrix}\square\\ \square\end{smallmatrix}}_b$	$\square\square_c$	$\overline{\square\square}_c$	$(\square_b,\square_d)$	$(\square_b,\square_c)$	$(\overline{\square}_b,\overline{\square}_d)$
1	3	3	3	0	3	1	2	0	3	1	1	0
2	3	3	3	0	3	1	2	3	0	1	1	0
3	3	3	0	3	3	0	3	0	3	1	0	1
4	3	3	0	3	3	0	3	3	0	1	0	1

Table 5: MSSM + 3 N_R spectrum for setups with $U(1)_Y = -\frac{1}{3}U(1)_a - \frac{1}{2}U(1)_b + U(1)_d$.

[4] In our displayed local D-brane configurations at least one of the left-handed quarks transforms as $(\square_a,\overline{\square}_b)$ under the D-brane gauge symmetry $U(3)_a \times U(2)_b$.

- The solution # 1: satisfies all constraints for matter parity, Baryon triality and hence also for Proton hexality. Matter parity R_2 and Baryon triality L_3R_3 are given by

$$R_2 = U(1)_a + U(1)_b + U(1)_c + 5U(1)_d \tag{4.1}$$
$$L_3R_3 = U(1)_a + U(1)_b + 3U(1)_c + U(1)_d \;. \tag{4.2}$$

Proton hexality takes the form

$$L_6^2R_6^5 = U(1)_a + U(1)_b + 9U(1)_c + 13U(1)_d \tag{4.3}$$

and does prevent the presence of R-parity violating couplings as well as the presence of dangerous dimension 5 proton decay operators, and at the same time allows for a μ-term as well as a Weinberg operator.

- The solution # 2: exhibits a massless $U(1)$ of the form

$$U^{add}(1) = U(1)_a + U(1)_b + U(1)_c - 3U(1)_d \tag{4.4}$$

which does not forbid any desired Yukawa couplings. The B-L symmetry is a linear combination of $U^{add}(1)$ and the hypercharge that takes the form $U(1)_{B-L} = 2U(1)_Y + \frac{1}{2}U^{add}(1)$. As before any discrete subgroup satisfies the constraints for the discrete symmetry (3.10) and (3.12). For instance the $\mathbf{Z}_2$ subgroup of U^{add} can be interpreted as matter parity. Moreover, one finds all four different discrete $\mathbf{Z}_3$ symmetries found in the MSSM using the pure field theoretical ansatz. They are given by the following linear combinations

$$L_3R_3^2 = 2U(1)_c + 4U(1)_d \tag{4.5}$$
$$L_3 = U(1)_a + U(1)_b + 5U(1)_c + 5U(1)_d \tag{4.6}$$
$$R_3 = U(1)_a + U(1)_b + U(1)_c + 3U(1)_d \tag{4.7}$$
$$L_3R_3 = U(1)_a + U(1)_b + 3U(1)_c + U(1)_d \;, \tag{4.8}$$

where R_3 originates from $U^{add}(1)$. Only Baryon triality L_3R_3 allows for the presence of a Weinberg operator. Thus in presence of the other discrete symmetries it is challenging to find a mechanism to generate neutrino masses. Finally, the setup also satisfies the constraints to exhibit all of the $\mathbf{Z}_6$ symmetries, i.e.

$$L_6^2R_6 = 3U(1)_a + 3U(1)_b + 19U(1)_c + 23U(1)_d \tag{4.9}$$
$$L_6^2R_6^3 = U(1)_a + U(1)_b + 17U(1)_c + 5U(1)_d \tag{4.10}$$
$$L_6^2R_6^5 = U(1)_a + U(1)_b + 9U(1)_c + 13U(1)_d \tag{4.11}$$
$$R_6 = U(1)_a + U(1)_b + U(1)_c + 21U(1)_d \;, \tag{4.12}$$

where R_6 originates from $U^{add}(1)$. In contrast to the solution # 1 here the proton hexality may be realized as a subgroup of a larger symmetry, namely a combination of the abelian gauge symmetry $U^{add}(1)$ and the discrete symmetry L_3R_3. In a concrete realization of this setup the $B \wedge F$ couplings may break the $U^{add}(1)$ down to matter parity R_2 and thus only Proton hexality survives in the low energy limit. In case a larger symmetry survives the

$B \wedge F$ couplings one needs a dynamical mechanism for the larger symmetry to break down to Proton hexality since otherwise the generation of a Weinberg operator and μ-term is not allowed.

- The solution # 3: may exhibit an additional $U^{add}(1) = U(1)_d$ which potentially remains massless, i.e. it satisfies the constraints (3.7). However, the presence of such an abelian gauge symmetry would spoil the model, since it would forbid various desired Yukawa couplings. Even worse there exists no discrete subgroup of the abelian gauge symmetry $U^{add}(1)$ that would allow the desired Yukawa couplings. Thus in a concrete realization it must be absent. The local D-brane configuration however does allow for a discrete $\mathbf{Z}_2$ that allows all desired Yukawa couplings, the matter parity R_2, given by
$$R_2 = U(1)_a + U(1)_b + U(1)_c + U(1)_d \tag{4.13}$$
which forbids all R-parity violating couplings.

- The solution # 4: may exhibit two additional $U(1)$'s given by
$$U_1^{add}(1) = U(1)_b - 2U(1)_c \qquad \text{and} \qquad U_2^{add}(1) = U(1)_d \ , \tag{4.14}$$
where the latter cannot survive as a gauge symmtery since it would forbid all desired Yukawa couplings. On the other hand the abelian gauge symmetry $U_1^{add}(1)$ does allow all superpotential terms. The B-L symmetry is given by $U_{B-L}(1) = \frac{1}{2}U_Y(1) - \frac{1}{4}U_1^{add}(1)$ in terms of the hypercharge and the additional $U_1^{add}(1)$. One finds for this configuration that the discrete subgroup of the two abelian gauge symmetries $U(1)_Y$ and $U_1^{add}(1)$ do give rise to matter parity R_2, to the $\mathbf{Z}_3$ symmetry R_3 and to the $\mathbf{Z}_6$ symmetry R_6 These discrete gauge symmetries are realized as the following linear combinations
$$R_2 = U(1)_a + U(1)_b + U(1)_c + U(1)_d \tag{4.15}$$
$$R_3 = U(1)_a + U(1)_b + U(1)_c \tag{4.16}$$
$$R_6 = U(1)_a + U(1)_b + U(1)_c + 9U(1)_d \ . \tag{4.17}$$
While R_2 forbids all R-parity violating couplings in this local D-brane configuration the absence of dimension 5 proton decay operators is rather accidental and does not originate from a discrete gauge symmetry.

4.2 Summary of the results

In [14] we performed an analogous bottom-up search for all hypercharge embeddings. Below we present the findings of that systematic search:

- The first thing to note is that we do not find in any of the semi-realistic D-brane configurations family dependent discrete gauge symmetries that allow for the desired Yukawa couplings $Q_L H_u U_R$, $Q_L H_d D_R$ and $L H_d E_R$. This is somewhat not expected since specifically the leptons in those D-brane configurations do arise from different intersections of D-brane stacks, and thus transform differently under the anomalous $U(1)$ factors. However, after determining the discrete charge of matter fields in different families we realise that they do have the same charges, even though their D-brane origin is significantly different. On the other hand, this allows us to compare the field theoretic prediction of CITE with our string theory results.

Disctere symmetries	Number of models
None	10
R_2	8
L_3R_3	8
R_2, R_3, R_6	3
R_2, L_3R_3, $L_6^2R_6^5$	1
R_2, L_3R_3, R_3, L_3, $L_3R_3^2$, $L_6^2R_6^5$, R_6, $L_6^2R_6^3$, $L_6^2R_6$	4

Table 6: Summary of the results. Ten models do not allow for any discrete gauge symmetry. The rest allow for some symmetries which are displayed on the left colon.

- In disagreement with the field theoretic result [6], we do not find any discrete $\mathbf{Z}_9$ and $\mathbf{Z}_{18}$ symmetries for the local MSSM D-brane configurations. This is due to the more constraining conditions for the appearance of discrete symmetries in D-brane compactifications.

- Matter parity R_2 is favoured for the hypercharge embeddings $U(1)_Y = -\frac{1}{3}U(1)_a - \frac{1}{2}U(1)_b$ and $U(1)_Y = -\frac{1}{3}U(1)_a - \frac{1}{2}U(1)_b + U(1)_d$ which appears for almost all D-brane setups with these hypercharge embeddings. For the hypercharge embedding $U(1)_Y = -\frac{1}{3}U(1)_a - \frac{1}{2}U(1)_b$ there is only one configuration out of 12 that allows for a $\mathbf{Z}_3$ and $\mathbf{Z}_6$ discrete symmetry. On the other hand for the hypercharge embedding $U(1)_Y = -\frac{1}{3}U(1)_a - \frac{1}{2}U(1)_b + U(1)_d$ we do find for each realization a $\mathbf{Z}_3$ symmetry, but only in two cases it is Baryon triality. Those local D-brane configurations also allow for Proton hexality.

  For the Madrid embedding $U(1)_Y = \frac{1}{6}U(1)_a + \frac{1}{2}U(1)_c - \frac{1}{2}U(1)_d$ almost all realizations have the potential to exhibit Baryon triality. However, the presence of matter parity is highly suppressed. Only for two setups we also find matter parity realized. Hence, those quivers pass the constraints to exhibit Proton hexality.

- Summarizing in table 6 we find only five local D-brane setups that have the potential to exhibit Proton hexality (out of 40), which is a particular intriguing discrete symmetry since it forbids all R-parity violating terms as well as all dangerous dimension 5 proton decay operators. This suggests that the presence of Proton hexality in D-brane compactifications is rather suppressed.

- Finally, one observes a similar pattern as in the field theoretical approach, namely that the presence of discrete $\mathbf{Z}_6$ symmetries is tied to the presence of $\mathbf{Z}_2$ and $\mathbf{Z}_3$ symmetries. We find the same relations as in pure field theory

$$R_2 \times L_3R_3 \cong L_6^2R_6^5\ , \quad R_2 \times R_3 \cong R_6\ , \quad R_2 \times L_3 \cong L_6^2R_6^3\ , \quad R_2 \times L_3R_3^2 \cong L_6^2R_6\ . \tag{4.18}$$

  Thus, the presence of R_2 along with a discrete $\mathbf{Z}_3$ symmetry implies the presence of a $\mathbf{Z}_6$ symmetry.

5. Conclusions

In [14], we study the presence of discrete gauge symmetries in D-brane compactifications. First, we translate the conditions for the presence of a discrete gauge symmetry in D-brane compactifications laid out in [12] into constraints on the transformation behaviour of the chiral matter fields. This allows for a bottom-up search, a search that does not require the knowledge of any features of the compactification manifold, for local D-brane configurations with respect to discrete gauge symmetries.

Next, we perform a systematic search for discrete gauge symmetries within a class of promising local D-brane quivers based on four stacks of D-branes. Those local configurations, that are consistent with the global consistency conditions, were found in [39, 50, 51] and exhibit the exact MSSM spectrum or the exact MSSM spectrum plus three right-handed neutrinos.

Within this class of intriguing four stack quivers there is no quiver that allows for a family dependent discrete gauge symmetry. Moreover, none of the local MSSM D-brane configurations exhibits a discrete $\mathbf{Z}_9$ and $\mathbf{Z}_{18}$ gauge symmetry, in contrast to the pure field theoretical approach [6]. Therefore, the string theory constrains on the transformation behaviour of the chiral matter fields for having a discrete gauge symmetry in D-brane compactifications goes beyond the four-dimensional discrete gauge anomaly conditions.

All $\mathbf{Z}_2$, $\mathbf{Z}_3$ and $\mathbf{Z}_6$ discrete gauge symmetries found in [6] can be also realized in the local D-brane configurations. We find that the realization of discrete symmetries depends on the hypercharge embedding of the D-brane configuration. For instance while the Madrid embedding favours Baryon triality it disfavours matter parity. The presence of Proton hexality, i.e. the simultaneous presence of matter parity and Baryon triality, is rather suppressed and only realized for five of the intriguing four D-brane-stack quivers. In those quivers the absence of R-parity and disastrous dimension 5 proton decay operators is not accidental, but can be explained by the presence of a discrete gauge symmetry.

It would be interesting to extend this analysis to local semi-realistic D-brane configurations with more than 4 D-brane stacks. Specifically, it would be interesting to see whether one can find family dependent discrete gauge symmetries in those realizations. Furthermore, another intriguing avenue is to extend the analysis to the NMSSM [53] and GUT realizations of the MSSM [40] as well as extending it to local D-brane configurations with additional exotics [54].

Finally, we would like to comment on the limits of the bottom-up approach applied here. The discrete gauge symmetries considered here purely originate from the anomalous $U(1)$ factors carried by each D-brane stack. In addition there may be abelian or even non-abelian gauge factors arising from isometries of the compactification manifold which can lead to abelian and non-abelian discrete gauge symmetries in the low energy effective action [13]. The consideration of discrete symmetries originating from isometries, requires the specification of the properties of the compactification and thus goes beyond the scope of this work.

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