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Branes with GUTs and Supersymmetry Breaking

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Abstract

We study Type I string theory compactified on a $T^6/\mathbb{Z}_3$ orientifold. The low-energy dynamics is most conveniently analyzed in terms of D3-branes. We show that a sector of the theory, which corresponds to placing an odd number of D3-branes at orientifold fixed points, can give rise to an $SU(5)$ gauge theory with three generations of chiral matter fields. The resulting model is not fully realistic, but the relative ease with which an adequate gauge group and matter content can be obtained is promising. The model is also of interest from the point of view of supersymmetry breaking. We show that, for fixed values of the closed string modes, the model breaks supersymmetry due to a conflict between a non-perturbatively generated superpotential and an anomalous $U(1)$ D-term potential.

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1. Introduction and Summary

The past few years have seen remarkable progress in our understanding of the non-perturbative behavior of string theory [1]. D-branes have played a vital role in these developments [2]. The consequences of this theoretical insight in string phenomenology are just beginning to be explored. In this paper we attempt to take a few preliminary steps in this direction. For related recent work see [3]–[8]. Perhaps the simplest idea to explore is that we live on a three dimensional brane or somewhat more precisely, that the $3+1$ dimensional spacetime corresponds to the world volume of a set of D3 branes. This immediately gives rise to a question: can a (grand unified) theory accommodating the standard model interactions and matter content be obtained in this manner?

D-brane model building is of interest from another point of view as well. Most of the model building so far has been carried out in the $E_8 \times E_8$ heterotic string [9]. In this context, there is a well known problem in reconciling, within the context of weakly coupled string theory, the "observed" unification of gauge coupling constants in supersymmetric extensions of the standard model [10] and the value of Newton's constant. Witten [3] has recently suggested working with the strongly coupled heterotic theory to avoid this problem. Another possibility, also mentioned in [3], is to consider model building in the Type I theory.

We begin this paper by considering, in Section 2, the question of gauge and gravity unification in the Type I string theory. We show that both the gauge coupling unification and the value of the Newton constant can be obtained within the context of Type I perturbation theory. Moreover, the analysis indicates that in several cases the more appropriate description is a T-dual one with D3 branes. This provides additional motivation to enquire about the standard model arising from D3 branes.

In Section 3, we turn to this issue by considering a compactification of the Type I theory on a $T^6/\mathbb{Z}_3$ orientifold. This compactification has been considered earlier by [11]. We point out that, in addition to the sector considered in [11], the moduli space for this compactification has additional disconnected branches, similar to the ones found in [12]. The different branches correspond to distinct ways in which the branes can be placed at the various orientifold fixed points. The additional branches of moduli space exhibit patterns of gauge symmetry breaking that are not otherwise allowed. In particular, we show in Section 3.1, that an $SU(5)$ grand unified theory with three generations of matter fields in the $\mathbf{10}$ and $\bar{\mathbf{5}}$ representations can arise in this manner. In Section 3.2, we show that some nonperturbative consistency conditions [13], [12] leading to the existence of these additional branches are met. The model obtained in this manner is not fully realistic: there are no Higgs fields present and there are Yukawa couplings violating baryon and lepton number. Even so, we view the relative ease with which an adequate gauge group and matter content can be obtained as encouraging.

Finally, in Section 4, we turn to another aspect of the $SU(5)$ theory mentioned above. The theory has an additional $U(1)$ gauge symmetry, which is anomalous. We show that a conflict between the non-perturbatively generated superpotential and the D-term of the anomalous $U(1)$ gives rise to supersymmetry breaking in this theory. Classically, the D-branes giving rise to the gauge theory are stuck at the orientifold plane. In the supersymmetry breaking vacuum, some of these branes are repelled by the orientifold and come to rest away from it. In this discussion of supersymmetry breaking, we neglect the gravitational interactions and keep the dilaton and a relevant orientifold blow-up mode fixed. We show that supersymmetry breaking occurs for any fixed values of the dilaton and orientifold blow-up mode. Once these modes are taken to be dynamical, there are, as usual, runaway directions along which supersymmetry is restored. What happens when the relaxation of the closed string modes and the gravitational interactions is included is an interesting question which we leave for the future.

2. Gauge Coupling Unification on D3 Branes

In this section, we discuss the constraints imposed on string model building by the requirement of gauge coupling unification (taking the values α_{GUT} and M_{GUT} for supersymmetric extensions of the standard model) and the observed value of Newton's constant. For the $E_8 \times E_8$ heterotic string, these requirements lead to the conclusion that string theory must be strongly coupled [3]. In contrast, as has been noted earlier in [3], we will see that in the case of the Type I string theory these requirements can be met while still working at weak string coupling. Moreover, the discussion below suggests that in several Type I models the six compactified dimensions can have a length somewhat bigger than the inverse GUT scale. In these cases, the gauge group and charged matter would arise from fields living on D3-branes that fill the $3+1$ dimensional flat spacetime.

The relation between the string scale α', Type I string coupling g_I, volume of compactification V_6, gauge coupling at unification $\alpha_{GUT} = g^2/4\pi$, and Newton's constant G_N is given by [1]:

$$G_N \;=\; \frac{(2\pi)^7}{16\pi}\,\frac{\alpha'^4}{V_6}\,g_I^2\ , \tag{2.1}$$

and

$$\alpha_{GUT} \;=\; \frac{(2\pi)^7}{4\pi}\,\frac{\alpha'^3}{V_6}\,g_I\ . \tag{2.2}$$

Here we will consider the situation where the six compactified dimensions have approximately the same size R. The volume V_6 is then roughly given by

$$V_6 \;=\; (2\pi R)^6 \tag{2.3}$$

To proceed, we need to decide how to relate the unification scale $M_{GUT} \simeq 10^{16}$ GeV to α' and R. In several string models the gauge couplings unify even in the absence of a grand unified group. In these cases, one expects the grand unification scale to correspond to the masses of the lightest extra charged states present in the string theory. These extra states can be of two kinds: Kaluza-Klein modes with a mass of order $1/R$, or higher string modes with a mass of order $1/\sqrt{\alpha'}$. If we assume that $R > \sqrt{\alpha'}$, the lightest extra states have a mass $m \sim 1/R$, leading to the relation $R \sim 1/M_{GUT}$. From eqs. (2.1), (2.2), and (2.3) it then follows that:

$$\frac{\alpha'}{R^2} = \frac{1}{\sqrt{2}} \alpha_{GUT} R M_{Pl}, \tag{2.4}$$

and

$$g_I = 4\sqrt{2} \frac{1}{\alpha_{GUT}^2 R^3 M_{Pl}^3}. \tag{2.5}$$

With the values $\alpha_{GUT} = 0.04$ and $R \sim 1/M_{GUT} = (10^{16}\text{GeV})^{-1}$ for the supersymmetric standard model [10], and $M_{Pl} = G_N^{-1/2} = 1.2 \times 10^{19}$ GeV, we get from eq. (2.5) that :

$$g_I \sim 10^{-6} . \tag{2.6}$$

Thus, the gauge coupling is small, as mentioned above. However from eq. (2.4) we find that:

$$\frac{\alpha'}{R^2} \sim 34 . \tag{2.7}$$

This shows that our starting assumption ($R > \sqrt{\alpha'}$) about the lightest extra charged states coming from Kaluza-Klein modes is incorrect. A consistent solution is obtained by assuming that $R < \sqrt{\alpha'}$. The lightest extra states which enter at the GUT scale are then higher string modes with mass $M \sim 1/\sqrt{\alpha'}$. In this case, the more appropriate geometrical picture is obtained by T-dualizing along the six compactified directions. Doing so turns the D9 branes into D3 branes. The T-dual radius and string coupling are given by:

$$\tilde{R} = \frac{\alpha'}{R} \tag{2.8}$$

and

$$\tilde{g}_I = \left(\frac{\alpha'}{R^2}\right)^3 g_I . \tag{2.9}$$

Eq. (2.2) then implies directly that

$$\tilde{g}_I = 2 \alpha_{GUT} \simeq 0.08 . \tag{2.10}$$

Furthermore, since the lightest excitations are higher string modes we now set $\alpha' \sim (M_{GUT})^{-2}$. Eq. (2.1) then gives

$$\tilde{R} = \sqrt{\alpha'} \left(\frac{1}{8} M_{Pl}^2 \alpha' \tilde{g}_I^2\right)^{\frac{1}{6}} \sim 3M_{GUT}^{-1} . \tag{2.11}$$

It is useful to describe the resulting picture in words. The gauge group and charged matter arises from D3-branes. The six compactified dimensions have a length scale somewhat bigger than the inverse GUT scale. In particular, we note that since all the degrees of freedom charged under the gauge groups arise from open strings that end on the three-branes, there are no momentum modes with mass of order $1/\tilde{R}$ charged under the gauge group. Instead, there are winding modes with a mass of order $\tilde{R}/\alpha'$ but these are somewhat heavier than the higher string modes with mass $\sim 1/\sqrt{\alpha'}$.

We should emphasize that the above picture is meant to be suggestive. Whether it applies or not will depend on the details of the compactification. It was noted in [3] that in the $E_8 \times E_8$ theory the large gauge coupling implies an extra dimension at a scale somewhat below the GUT scale (for recent work and a list of references, see [4]). Here it is interesting to note that the presence of extra large dimensions might be true in the Type I case as well, and more generally, in attempts to build string models involving branes. We should also note that the conclusion with regards to the smallness of the string coupling is secure regardless of the exact relation between R and $\sqrt{\alpha'}$.[2]

We end this section with a few comments. First, strictly speaking, the discussion above applies to models where gauge coupling unification occurs in the absence of a grand unified group. One can ask what happens if the low-energy field theory is a grand unified theory. In this case the lightest extra string states need not occur at the GUT scale but could have larger masses. Eq. (2.1) then shows that in these cases the compactification scale $\tilde{R}$ should be comparable to α'. Even so, as we see in the next section it might be sometimes convenient to analyze such a model in terms of D3-branes. Second, in our analysis we have taken all the compact dimensions to have roughly the same size. This of course need not be true. For recent discussions of large extra dimensions and weak scale strings see [5]. Finally, we have assumed that the gauge group and matter content arises from the perturbative sector of Type I theory. This, too, need not be true. One could have a situation where some of the degrees of freedom arise from 9-branes while others arise from 5-branes; for model building along these lines see [7], [8], [14].

3. A Three Generation Model on D3 Branes

In this section, we present a simple "three generation" model with D3 branes placed at a $T^6/\mathbb{Z}_3$ orientifold. The model is, admittedly, not a realistic one, but it will serve the purpose of making several generic points quite explicit. In general, the moduli space of the gauge theory which governs the low-energy dynamics can be quite complicated with several disconnected sectors. The D3-brane picture allows for a geometric description of these different branches of moduli space [12]. The different branches correspond to the distinct ways in which the branes can be placed at the various orientifold fixed points.

The additional branches of moduli space can have multiple uses. We will see below that

[2]For example, setting $\sqrt{\alpha'} = R$ in eq. (2.2) still gives (2.10) for the gauge coupling.

they exhibit interesting patterns of gauge symmetry breaking that are not otherwise possible. In addition, branes placed at different orbifold fixed points can serve as "visible" and "hidden" sectors; the latter can be responsible for supersymmetry breaking. The lightest excitations of strings stretching between branes at different fixed points transform as fundamentals under both the "hidden" and "visible" gauge groups; these could be instrumental in communicating supersymmetry breaking.

3.1 The $\mathrm{T}^6/\mathbb{Z}_3$ *orientifold*

We now turn to studying the Type I theory compactified on a $\mathrm{T}^6/\mathbb{Z}_3$ orientifold. This theory has been analyzed by [11] and more recently by [14], where the low energy dynamics was shown to correspond to an $SU(12) \times SO(8)$ gauge theory. Our main purpose here will be to study some sectors of moduli space which are disconnected from the $SU(12) \times SO(8)$ theory mentioned above. For this purpose it will be often convenient to T-dualize the Type I theory along the six directions of T^6. Doing so turns the 9-branes into 3-branes. The disconnected sectors then correspond to placing an odd number of 3-branes at the orientifold fixed planes and can be easily visualized. We show below how an $SU(5)$ theory with three generations of matter fields in the $\mathbf{10}$ and $\bar{\mathbf{5}}$ representations can be obtained in this manner.

Let us describe the $\mathrm{T}^6/\mathbb{Z}_3$ orientifold in more detail. We work for the most part in the T-dual description involving D3 branes. The D3 branes stretch along $X^{1,2,3}$. We introduce complex coordinates, $z_1 = X^4 + iX^5, z_2 = X^6 + iX^7, z_3 = X^8 + iX^9$, in the compactified six-dimensional space. Consider the two-torus obtained by identifying points under

$$z \simeq z + R \simeq z + R\, e^{\frac{i2\pi}{3}} . \tag{3.1}$$

The T^6 is obtained by taking three copies of this two torus, corresponding to the three complex coordinates z_1, z_2, z_3. The orientifold group is given by:

$$G = \{ 1,\ \alpha,\ \alpha^2,\ \Omega R(-1)^{F_L},\ \Omega R(-1)^{F_L}\alpha,\ \Omega R(-1)^{F_L}\alpha^2 \} . \tag{3.2}$$

Here, α is a spacetime symmetry whose action is given by:

$$(z_1,\ z_2,\ z_3) \rightarrow (\alpha z_1,\ \alpha z_2,\ \alpha z_3) . \tag{3.3}$$

Ω denotes world-sheet orientation reversal, and R is a reflection $z_i \rightarrow -z_i, i = 1, 2, 3$. F_L is an operator that flips the sign of the left-moving Ramond states. The orientifold group G has a $\mathbb{Z}_2$ subgroup

$$G_{\mathbb{Z}_2} = \{1, \Omega R(-1)^{F_L}\} \tag{3.4}$$

and a $\mathbb{Z}_3$ subgroup

$$G_{\mathbb{Z}_3} = \{1, \alpha, \alpha^2\} . \tag{3.5}$$

These will play a useful role in the subsequent discussion.

In addition to acting on the spacetime indices, the orientifold group acts on the Chan-Paton indices of the open string states stretching between D3 branes [2]. The action of the group elements, $\Omega R(-1)^{F_L}$ and α, on the Chan-Paton factors λ is:

$$\lambda \to \gamma_{\Omega R(-1)^{F_L}} \ \lambda^T \ \gamma^{-1}_{\Omega R(-1)^{F_L}} \ , \tag{3.6}$$

and

$$\lambda \to \gamma_\alpha \ \lambda \ \gamma_\alpha^{-1}. \tag{3.7}$$

The matrices γ_α and $\gamma_{\Omega R(-1)^{F_L}}$ must furnish a representation of the orientifold group. The matrices $\gamma_{\Omega R(-1)^{F_L}}$, representing the action of the $\mathbb{Z}_2$ part of the orientifold group should obey [2]:

$$\gamma_{\Omega R(-1)^{F_L}} \ = \left(\gamma_{\Omega R(-1)^{F_L}}\right)^T \ . \tag{3.8}$$

In the absence of the $\mathbb{Z}_3$ orbifold projection, the $\Omega R(-1)^{F_L}$ projection would lead to an SO gauge group on the D3 brane world volume.

Tadpole cancellation conditions play an important role in ensuring the consistency of the string compactification. For the $\mathrm{T}^6/\mathbb{Z}_3$ orientifold these were discussed in [11]. For the sake of brevity we will not discuss a detailed derivation of these conditions here. Instead we will content ourselves with stating them; as the reader will see these conditions give rise to anomaly free gauge theories.

As expected, the untwisted Ramond-Ramond 4-form charge conservation conditions require the presence of 32 D3 branes to cancel the orientifold charge. In addition, there are charge cancellation conditions for the twisted RR fields. Before stating these, it is useful to consider the action of the $G_{\mathbb{Z}_3}$ and $G_{\mathbb{Z}_2}$ subgroups of the orientifold group on the T^6. Consider first a two torus shown in Fig. 1.

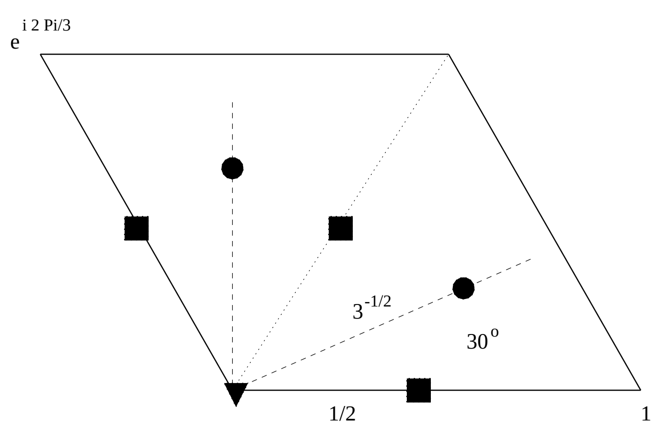

Figure 1

The origin, denoted by the triangle in Fig. 1, is the only fixed point with respect to the full $\mathbb{Z}_6$ orientifold group. In addition the $G_{\mathbb{Z}_3}$ subgroup has two additional fixed points, at $z = \frac{1}{\sqrt{3}}e^{\frac{i\pi}{6}}$ and $z = \frac{1}{\sqrt{3}}e^{\frac{i\pi}{2}}$, respectively. These two fixed points are interchanged under the action of the $G_{\mathbb{Z}_2}$ subgroup and are denoted by circles in Fig. 1. Similarly, the $G_{\mathbb{Z}_2}$ subgroup has three additional fixed points at $z = \frac{1}{2}, \frac{1}{2}e^{\frac{i2\pi}{6}}, \frac{1}{2}e^{\frac{i2\pi}{3}}$, which are denoted by squares. These are images of each other under the $G_{\mathbb{Z}_3}$ symmetry. The fixed points for the T^6 can now be deduced in a straightforward manner. There is only one fixed point under the full $\mathbb{Z}_6$ symmetry—the origin. In addition, as in the T^2 case there are fixed points of the $G_{\mathbb{Z}_3}$ symmetry which are transformed into one another by the $G_{\mathbb{Z}_2}$ and vice-versa; these will be referred to, in what follows, as $G_{\mathbb{Z}_3}$ and $G_{\mathbb{Z}_2}$ fixed points, respectively.

We can now return to the tadpole conditions for the twisted Ramond-Ramond fields. If γ_α is the matrix which represents the action of $G_{\mathbb{Z}_3}$ on the branes at the origin, one finds [11] that:

$$\mathrm{Tr}\ \gamma_\alpha \ = \ -4\ . \tag{3.9}$$

In contrast, at a $G_{\mathbb{Z}_3}$ fixed point one finds that

$$\mathrm{Tr}\ \gamma_\alpha \ = \ 0. \tag{3.10}$$

Finally, at each $G_{\mathbb{Z}_2}$ fixed point (and consistently at its $\mathbb{Z}_3$ image points) one can choose to place an even or odd number of branes.

The simplest way to meet these conditions is to place all the 32 D3 branes at the origin. This gives rise, from eq. (3.9), to a gauge theory with $SU(12) \times SO(8)$ gauge group with three generations of $(\overline{\square},\square) + (\begin{smallmatrix}\square\\\square\end{smallmatrix},\mathbf{1})$ fields which was discussed in [11].

The rank of the $SU(12) \times SO(8)$ gauge symmetry can be reduced by moving some of the branes away from the origin in a continuous manner. To be consistent with the $\mathbb{Z}_6$ orientifold symmetry, however, these branes can only be moved away from the origin in sets of six. From eq. (3.9) it then follows that the rank of the $SU(N)$ factor must always be odd; this precludes an $SU(5)$ gauge symmetry which is attractive from a phenomenological point of view.

3.2 The $SU(5)$ theory

We turn now to exploring some branches of moduli space, which are disconnected from the $SU(12) \times SO(8)$ theory mentioned above. We will see how some of these branches give rise to an $SU(5)$ gauge theory with three generations of fields in the $\mathbf{10}$ and $\bar{\mathbf{5}}$ representations. We discuss examples of such disconnected branches below, but before doing so it is worth summarizing the essential features responsible for the grand unified theory.

In some sectors of moduli space, an odd number of branes can be removed from the origin. In particular, a situation can arise where only 11 of the 32 D3 branes are left at the origin. Eq. (3.9) then implies that one has an $SU(5)$ gauge symmetry. In addition there is an anomalous $U(1)$ symmetry (which is broken at the string scale). The theory has $N = 1$ supersymmetry with matter content corresponding to three generations of matter fields which transform as:

	$SU(5)$	$U(1)$
$A_{i=1,2,3}$	$\begin{smallmatrix}\square\\\square\end{smallmatrix}$	2
$\bar{Q}_{i=1,2,3}$	$\bar{\square}$	-1

. (3.11)

The theory has a renormalizable tree-level superpotential given by

$$W_{tree} = \epsilon^{ijk} A_i \bar{Q}_j \bar{Q}_k. \tag{3.12}$$

The three generations arise because the $\mathbb{Z}_3$ action in eq. (3.3) does not distinguish between the three (complex) transverse coordinates, thus one set of matter fields in eq. (3.11) arise from each of them.

We now turn to discussing how these disconnected branches arise. As we saw in the discussion above, if we start with all 32 branes at the origin and move some away in a continuous manner, one is always left with an even number of branes at the origin. Thus to get 11 branes some of them must be placed at fixed points of the $G_{\mathbb{Z}_3}$ of $G_{\mathbb{Z}_2}$ subgroups, eqs. (3.4), (3.5). Now eq. (3.10) implies that the number of branes at a $G_{\mathbb{Z}_3}$ fixed point must be a multiple of three. In addition, as we saw above, each $G_{\mathbb{Z}_3}$ fixed point has a $G_{\mathbb{Z}_2}$ image. Thus, one finds that all the branes at a $G_{\mathbb{Z}_3}$ fixed point can be moved continuously away in a $\mathbb{Z}_6$ symmetric manner back to the origin. Disconnected branches of moduli space can be obtained, however, by placing an odd number of branes at a $G_{\mathbb{Z}_2}$ fixed point (and its two images under $G_{\mathbb{Z}_3}$). Since there are a large number of $G_{\mathbb{Z}_2}$ fixed points in T^6 this gives rise to a large number of possibilities. We will not analyze all of them in detail here. Rather, as an illustrative example, we focus on a case that gives rise to the $SU(5)$ "grand unified theory" mentioned above.

For this purpose, the simplest possibility is to consider a situation where all the branes are at the origin as far as the third T^2 (corresponding to the z_3 coordinate) is concerned, but not as far as the other two tori are concerned. Consider placing one D3 brane at a $G_{\mathbb{Z}_2}$ fixed point in the first T^2 and at a $G_{\mathbb{Z}_2}$ fixed point in the second T^2—this brane has two images under the $\mathbb{Z}_3$ symmetry. Next, place one D3 brane at the origin of the first two-torus, but at a $G_{\mathbb{Z}_2}$ fixed point of the second T^2—this brane has two images as well. Finally, place a D3 brane at a $G_{\mathbb{Z}_2}$ fixed point of the first torus, but at the origin of the second. This brane has two images as well. Altogether, counting images, this gives us 9 D3 branes—an odd number—which are stuck to orientifold planes away from the origin (note, however, that the number of D3 branes at the $G_{\mathbb{Z}_2}$ fixed points in each of the three two-tori is even; this is required by the consistency conditions discussed in the following section). The remaining

23 branes at the origin give rise to an $SU(9) \times SO(5)$ gauge symmetry. The matter fields transform as three generations of $(\begin{smallmatrix}\square\\\square\end{smallmatrix}, \mathbf{1})$ and $(\overline{\square}, \square)$, under the $SU(9) \times SO(5)$ gauge group.[3]

Finally, one can move 12—two sets of six—of the remaining branes away from the origin, leaving behind 11 branes, obtaining thus the $SU(5)$ theory mentioned above, eq. (3.11). We also should mention that for a generic position of these 12 branes, the full gauge symmetry includes an additional $U(1) \times U(1)$ factor. This "hidden sector" gauge group can be further enhanced if the branes are placed at $G_{\mathbb{Z}_3}$ fixed points or at the $G_{\mathbb{Z}_2}$ orientifold planes. For example, placing six of the 12 branes at a G_{Z_3} fixed point (and the remaining six at the image point) gives rise to an $N = 1$ theory with $SU(2)^3$ symmetry and three sets of chiral matter transforming as bifundamentals under pairs of the $SU(2)$'s. Dividing the 12 branes between a $G_{\mathbb{Z}_2}$ orientifold plane and its two images, on the other hand, can give rise to a theory with $N = 4$ supersymmetry and an $SO(4)$ or $SO(5)$ gauge symmetry (the $SO(5)$ symmetry can arise if the orientifold planes chosen already contain a D3 brane stuck to them, as mentioned above in the discussion of the disconnected moduli space).

So far, we have ignored the effects of open strings stretched between branes at different fixed points. The lightest excitations of such strings are massive states which transform as fundamental-antifundamental under the respective world volume gauge groups. In the example we gave above, there can be two world volume theories with $N = 1$ supersymmetry. If supersymmetry were dynamically broken in one of these theories, supersymmetry breaking would be communicated to the other gauge theory via the massive chiral multiplets just described (and, of course, by the supergravity in the bulk). A more precise investigation of this would probably involve details of the supersymmetry breaking dynamics and the stabilization of the dilaton [15]; we leave this for future investigation.

3.3 Non-perturbative consistency conditions

There is one subtlety concerning disconnected sectors of moduli space that needs to be mentioned. Sometimes such sectors are not allowed, even when they pass all the perturbative consistency conditions, due to non-perturbative reasons. Similar issues were addressed in [13], [12]. We will not be able to discuss this matter in full detail here, but will mention some salient points. The basic idea behind the non-perturbative consistency conditions is as follows. The Type I theory does not have any perturbative states which transform in spinor representations of $SO(32)$. However, such states are present in the dual heterotic $SO(32)$ theory and are nonperturbative in the Type I theory. Allowing for such spinor representations imposes additional consistency conditions—whose origin from the Type I viewpoint is non-perturbative.

We can verify that the example discussed above, giving rise to the $SU(5)$ model, meets

[3] We are using a notation where the vector of $SO(N)$ is denoted by $\square$.

various non-perturbative conditions. However, we should caution the reader that there might be other conditions, besides the ones we have checked that might not be met.[4] It is in fact useful to carry out this discussion in the original description of Type-I theory in terms of 9 branes. The essential feature giving rise to the disconnected branch of moduli space in the example above was the fact that there were 9 D3 branes (counting images) which were "stuck" at the orientifold plane. In the T-dual 9-brane language we are using now, the positions of branes correspond to expectation values for particular Wilson lines. The question is whether the Wilson lines' expectation values are consistent with the existence of states that transform as $SO(32)$ spinors—the holonomies around any contractable loop should be trivial in the appropriate spinor representation.

The example discussed in Section 3.2 is equivalent to turning on four Wilson lines along the noncontractible loops of two of the two-tori. It turns out in this case that the holonomy around any contractible path is trivial in all representations of $SO(32)$. One can show this by explicitly writing down the Wilson lines that correspond to the brane configuration with 9 D3 branes removed from the origin, which was described in Section 3.2. Moreover, in this example, an explicit periodic flat connection, which is not constant on the z_1, z_2 four-torus, can be found. This can be done by a straightforward generalization of the construction of ref. [16] to the case of T^4. Furthermore, as in [13], one can show that if W is the Wilson line relevant for the particular $\mathbb{Z}_3$ fixed point then $(W\gamma_\alpha)^3 = 1$ in spinor representations as well.

Before moving on, let us mention that the example of Section 3.1, giving rise to 11 branes at the origin, is just one of several possibilities consistent with the various conditions. For example, one can easily work out brane configurations with an odd total number of branes removed from the origin that involve moving branes to the $G_{\mathbb{Z}_2}$ fixed points in all three two-tori.

3.4 The GUT: shortcomings

It is useful to describe the construction of the $SU(5)$ theory in group theoretic terms perhaps more familiar to some model builders. The $SO(32)$ gauge symmetry is broken to an $SO(11)$ subgroup (times a hidden sector group). The orientifold projection then further breaks the symmetry to $SU(5)$ (with an additional anomalous U(1)). The $\bar{\mathbf{5}}$ and $\mathbf{10}$ matter fields arise from the adjoint representation of $SO(11)$ by the orientifold projection. The three generations arise because there are three complex (six real) transverse dimensions and because the orientifold group acts in an identical manner on the three directions.

The relative ease with which a realistic gauge group and matter content can be obtained in the Type I theory is interesting. We should note, that even though we used a D3 brane

[4]By way of comparison, we note that the conditions we have tested are the analogue of those discussed in Section 2.3 of [13]. The authors also formulated a stronger set of conditions in Section 5 of ref. [13]. We have not investigated the presence of such stronger consistency conditions in the present case.

description to simplify the discussion, the construction as such was purely in the context of perturbative Type I string theory. In particular, the matter content was obtained, even though we did not have any spinor representations of $SO(32)$ to begin with—in fact all the matter fields can be thought of as being obtained by truncating adjoint representations of $SO(32)$.

However, it should also be noted that the model is meant as an illustrative example and is not realistic. There are several reasons for this. First, there are no Higgs fields either to break the $SU(5)$ gauge symmetry or to give rise to the $SU(2)$ Higgs doublets of the supersymmetric standard model. Second, and this is perhaps a more important limitation, as was mentioned in passing in eq. (3.12) above there is a Yukawa coupling in the theory that violates baryon and lepton number. The underlying reason for this coupling is that in the $N = 4$ theory, which can be thought of as the starting point for the above construction, there is a coupling involving the three adjoint fields corresponding to the transverse directions. In the case of, say, spinor representations of $SO(10)$, a trilinear $\mathbf{16}^3$ coupling is not allowed by gauge invariance and an R-parity symmetry can often be imposed to prevent baryon- and lepton-number violating terms. However, in the present example, where all the matter arises from adjoint representations no such R parity symmetry is present. This limitation is likely to be quite general.

4. Supersymmetry Breaking

We turn now to another feature of the $SU(5) \times U(1)$ theory discussed above. As we will see below, in the world-volume field theory context, a conflict between the non-perturbatively generated superpotential and the anomalous $U(1)$ D-term results in the breaking of supersymmetry in this theory. Our discussion of supersymmetry breaking will only involve the open string sector corresponding to the world volume theory on the D3 branes. Gravity and other closed string effects will be neglected. In particular, the dilaton and the orientifold blow-up mode [19], which acts as the Fayet-Illiopoulos term for the $U(1)$, are regarded as coupling constants, and will be kept fixed in the discussion below. We will establish that supersymmetry breaking occurs for any finite value of these couplings.[5] But, as is usually the case, once they are allowed to vary, we find that there are runaway directions along which supersymmetry is restored. Perhaps, these could be stabilized by (yet poorly understood) nonperturbative corrections to the Kähler potential (see, e.g. the recent discussion in [15] and references therein). The stability of the supersymmetry breaking ground state in the context of the full theory is a complicated issue, about which we have nothing to say here.

Towards the end of this section, we will briefly comment on this runaway behavior and

[5] More accurately, supersymmetry breaking will be shown when the string coupling is small enough to argue with confidence that the low-energy dynamics is governed by the $SU(5) \times U(1)$ theory.

the gravitational back reaction. One final comment before we get started: the discussion in this section only involves the branes at the $\mathbb{Z}_6$ fixed point, any details of compactification etc. are irrelevant in this context. For example, the analysis here applies to the world volume theory of 11 D3 branes on the noncompact $\mathbb{C}^3/\mathbb{Z}_3$ orientifold as well.

4.1 Supersymmetry breaking in the $U(5)$ theory

Our strategy to establish supersymmetry breaking is as follows. We first neglect the anomalous $U(1)$ and show that in the $SU(5)$ theory the resulting non-perturbative superpotential gives rise to runaway behavior. Then on incorporating the anomalous $U(1)$ we find that its D-term gives rise to an energy that grows along the runaway directions. This leads to supersymmetry breaking.

The $SU(5)$ "three generation" model is an s-confining theory [17]. The infrared degrees of freedom are the mesons and baryons

$$\begin{aligned} C &= A\cdot\bar{Q}\cdot\bar{Q} &&\sim (\mathbf{3},\ \bar{\mathbf{3}}\ ,0)\ , \\ B &= A^5 &&\sim (\mathbf{6},\ \mathbf{1}\ ,10)\ , \\ M &= A^3\cdot\bar{Q} &&\sim (\mathbf{8},\ \mathbf{3}\ ,5)\ , \end{aligned} \tag{4.1}$$

where we have shown their transformation properties under the global $SU(3)_A \times SU(3)_{\bar{Q}}$ symmetry and the last column in each entry refers to the charges under the anomalous $U(1)$ which follow from eq. (3.11). The confining superpotential is [17]:

$$W = \frac{C_a^\alpha\ B^{\beta\gamma}\ M_\gamma^{\delta a}\ \epsilon_{\alpha\beta\delta} + M_\beta^{\alpha a}\ M_\gamma^{\beta b}\ M_\alpha^{\gamma c}\ \epsilon_{abc}}{\Lambda^9} + \lambda\ \delta_\alpha^a\ C_a^\alpha\ , \tag{4.2}$$

where $a, b, ...(\alpha, \beta, ...)$ denote indices under the $SU(3)_{\bar{Q}(A)}$ symmetry, respectively, and the last term is the tree-level superpotential. The tree-level superpotential breaks the global symmetry to the diagonal $SU(3)_{diag}$. It lifts all the C flat directions, but does not lift the B and some of the M directions. The superpotential coupling λ in (4.2) is proportional to the value of the gauge coupling at the string scale (since the tree-level superpotential is the projection of the $N = 4$ superpotential).

We will show now that the F-term equations of motion following from (4.2) have no solutions for finite field expectation values. Consider the equations of motion following from the superpotential (4.2) (suppressing numerical constants):

$$M_\alpha^{\beta a}\ B^{\alpha\lambda}\ \epsilon_{\beta\lambda\delta} = \delta_\delta^a\ , \tag{4.3}$$

$$\epsilon_{abc}\ M_\alpha^{\beta b}\ M_\beta^{\gamma c} + \epsilon_{\alpha\beta\delta}\ C_a^\beta\ B^{\delta\gamma} = 0\ , \tag{4.4}$$

$$\epsilon_{\gamma\alpha\beta}\ C_a^\alpha\ M_\delta^{\beta a} = 0\ . \tag{4.5}$$

Multiplying the first equation (4.3) by $\epsilon^{\delta\mu\nu} M_\nu^{\gamma d} \epsilon_{bad}$, summing over δ and a, and substituting for $\epsilon_{abc} M_\alpha^{\beta b} M_\beta^{\gamma c}$ from (4.4), we obtain:

$$- \ \epsilon_{\alpha\delta\beta} \ B^{\alpha\mu} \ B^{\delta\gamma} \ C_b^\beta \ + \ \epsilon_{bad} \ M^{\mu a} \ B^{\alpha\nu} \ M_\nu^{\gamma d} \ = \ M_b^{\gamma\mu} \ - \ \delta_b^\mu \ M_\nu^{\gamma\nu} \ . \tag{4.6}$$

Now under $SU(3)_{diag}$, the field $M_\beta^{\alpha a}$ decomposes as a **3**, given by the partial trace $M_a^{\alpha a}$, a $\bar{\mathbf{6}}$ which is antisymmetric in the upper two indices, and a **15** which is symmetric in the upper indices (and traceless). Note that the l.h.s. of (4.6) is antisymmetric in μ, γ, hence only the r.h.s. contributes to the symmetric part of M. This gives rise to the relation:

$$M_b^{\gamma\mu} \ + \ M_b^{\mu\gamma} \ - \ \delta_b^\mu \ M_\nu^{\gamma\nu} \ - \ \delta_b^\gamma \ M_\nu^{\mu\nu} \ = \ 0 \ , \tag{4.7}$$

from which in turn it follows that $M_a^{\alpha a}$ vanishes as does the symmetric part of M. Thus the **3** and **15** components of M are zero. The remaining $\bar{\mathbf{6}}$ can be written as

$$M_a^{\gamma\mu} \ = \ \epsilon^{\gamma\mu\kappa} \ s_{\kappa a} \ , \tag{4.8}$$

where s is symmetric in the two indices. Substituting into eq. (4.3), and evaluating for $B \cdot s$ then leads to the relation:

$$s \ = \ - \frac{1}{2} \ B^{-1} \ . \tag{4.9}$$

Substituting eq. (4.8) in (4.5) and noting that s is invertible, eq. (4.9) then leads to $C_a^\alpha = \delta_a^\alpha \mathrm{Tr}\ C$. This then implies that $C = 0$. Finally, substituting into eq. (4.4), one similarly finds that $s_{\alpha\beta} = \delta_{\alpha\beta}$ Tr s. This leads to the conclusion that $s = 0$. But now we see from eq. (4.9) that B must go to infinity. Thus we have established that there are no solutions to the F flatness conditions at finite expectation values.

The equations (4.3)–(4.5) do have runaway solutions. The discussion above leads to the conclusion that along a runaway direction, C and $M \to 0$, while $B \to \infty$ in an invertible manner—more precisely $B^{-1} \to 0$. For example, a runaway vacuum solution with $SU(3)_{diag} \to SO(3)$ global symmetry is $B^{\alpha\beta} \sim \delta^{\alpha\beta} b$, $C^{\alpha a} \sim \delta^{\alpha a} b^{-3}$, $M^{\alpha a\beta} \sim \epsilon^{\alpha a\beta} b^{-1}$, with $b \to \infty$. The physics along the B, det $B \neq 0$ directions is easy to understand. Along these directions the mesons C and M obtain mass. Upon integrating them out, the superpotential of the low-energy theory is $W_{eff} = \lambda^3 \Lambda^{18}/\mathrm{det} B$, showing explicitly the runaway behavior.

So far we have studied the non-perturbative superpotential. Now let us include the anomalous $U(1)$ by "turning it on" in the effective theory of the mesons C, M, and B. The last column in each row of eq. (4.1) gives the $U(1)$ charges of the three fields. We see that C has charge zero, while B and M both have positive charge. We have argued above that along a runaway direction B must go to infinity (in an invertible manner) while M and C go to zero. Since all components of B have positive charge with respect to the $U(1)$, we see that along such a direction the $U(1)$ D-term contribution to the energy blows up. Thus, the runaway behavior dictated by the non-perturbative superpotential is in conflict with the

$U(1)$ D-term potential, leading to the breaking of supersymmetry (note that this is similar to the mechanism of ref. [18]).

4.2 Remarks on supersymmetry breaking

We end this discussion of supersymmetry breaking with a few remarks. We first remind the reader of an important feature of type I orbifolds: the anomaly cancellation of the "anomalous" $U(1)$s is provided, as in the heterotic case, by a Green-Schwarz mechanism. However, unlike the heterotic case, the axion that shifts under the $U(1)$ to cancel the anomaly is a model-dependent field—the twisted Ramond-Ramond field from the closed string sector (this has been pointed out in [19] and recently discussed in [20]). It is in the same supermultiplet as the orbifold blow-up mode (the twisted NS-NS field) and can be described in terms of a chiral superfield, denoted hereafter by C, with a kinetic term

$$\int d^4\theta \, \left(C \, + \, C^\dagger \, + V\right)^2 \, + \, \ldots \, , \tag{4.10}$$

where dots denote higher-order terms. The leading term (4.10) can be written by demanding $U(1)$ invariance and a smooth kinetic term for C in the orbifold limit $\langle C \rangle = 0$. Here V is the anomalous $U(1)$ vector superfield; in addition to (4.10), the field C also has a Wess-Zumino coupling to the gauge field strengths, of the form $\int d^2\theta C W^\alpha W_\alpha$ [19]. In a superunitary gauge, the term (4.10) represents a mass term (of order the string scale) for the anomalous $U(1)$ vector superfield. By giving an expectation value to the real part of C (blowing up the orbifold) one can induce "tree-level" FI terms, with $\zeta_{FI}^2 \sim \langle C + C^\dagger \rangle$, as follows from (4.10). That (4.10) is correct follows from the computation of ref. [19] of the coupling of the real part of C (the twisted NS-NS field) to the D-term of the vector superfield (and from a subsequent supersymmetry transformation). This coupling arises from the disk with two scalar vertex operators attached to the boundary, and a closed string twisted NS-NS scalar vertex operator in the bulk [19], and is of order $\mathcal{O}(g_{string}) \sim g^2$.

We note that the conclusion regarding supersymmetry breaking is true for any sign and finite value of the $U(1)$ Fayet-Illiopoulous term. Depending on the sign of the FI term, the D-term potential can have a zero at finite values of the fields.[6] However the F-term potential vanishes only at infinity. Thus supersymmetry is broken. It is possible though to have vanishing D- and F-term potentials for infinite (negative) value of the FI term.

One would like to find out where the resulting supersymmetry breaking vacuum lies. Unfortunately, this is quite difficult—as the following argument shows, one expects the vacuum to lie in a strongly coupled region where a semiclassical analysis is not applicable. Assuming first that such an analysis is valid, upon balancing the U(1) D-term energy with

[6]Since the $U(1)$ is anomalous one expects that the FI term is renormalized at one loop. However, an explicit calculation shows that the FI term is not generated in open string orbifolds, because of cancellations between contributions of worldsheets of different topology [21].

the F-term energy one finds that the vacuum energy (with vanishing FI term) scales as $g_1^{16/9}\lambda^{2/3}\Lambda^4$, while the typical expectation value of a field goes like $v \sim \lambda^{1/6} g_1^{-1/18}\Lambda$. Here g_1 is the $U(1)$ gauge coupling, λ is the tree-level Yukawa coupling, eq. (4.2), and Λ is the strong coupling scale of the $SU(5)$ gauge theory. For a semiclassical analysis to be valid, one requires $v \gg \Lambda$. If the gauge coupling g_1 and λ were independent parameters, this could have been achieved by taking g_1 to zero keeping λ fixed. However, in our case $g_1 \sim \lambda$, thus for small g_1 the vacuum lies in the strongly coupled region and the semi-classical analysis is not applicable. One could make $v \gg \Lambda$ by taking $g_1 \gg 1$, but then the string coupling would be large, again making a semiclassical analysis invalid.

While we cannot determine the vacuum explicitly, we know that some of the B, M, and C fields must get vacuum expectation values. These expectation values should correspond to displacing some of the D3 branes away from the orientifold. We remind the reader that the $SU(5)$ theory under consideration here is the world volume theory for 11 D3 branes placed at the orientifold. Classically, the 11 D3 branes are all stuck to the orientifold plane and the configuration has no moduli. This corresponds to the fact that in order to meet the tadpole conditions and respect the $\mathbb{Z}_6$ symmetry no branes can be moved away from the orientifold point. Quantum mechanically, due to non-perturbative supersymmetry breaking effects we see that some of the branes are repelled by the orientifold plane and come to rest away from it so as to minimize the energy. Since there are no moduli the configuration of D3 branes cannot be described in terms of classical geometry. The displacement of branes which are classically stuck at the orientifold is somewhat reminiscent of the splitting of orientifold 7-planes discussed in ref. [22].

As was mentioned in the beginning of this section, the above analysis neglected all interactions with closed string sector modes. One might at first expect that gravitational interactions are small at low-energies and so can be neglected. But once supersymmetry is broken and a (boundary) cosmological constant is induced, this is not apriori true. Also, interactions with some other closed string modes, which determine the couplings of the brane theory are important. There are two modes of this kind. A blow-up mode for the orientifold determines the FI term of the $U(1)$ [19]; this mode (together with its partner) also enters in the determination of the coupling constant (theta angle) for the $SU(5)$ theory. Similarly the dilaton determines the gauge coupling of the $U(1)$ and together with the blow-up mode mentioned above determines the $SU(5)$ gauge coupling. As we have argued here supersymmetry breaking occurs for any fixed values of these couplings, but there are runaway directions along which it can be restored. For example, as we saw above there is a direction along which the FI term can go to infinity with appropriate sign. Similarly, if the dilaton goes to infinity, supersymmetry is restored. In fact, as has been argued recently in [15], it is necessary to include the dynamics of these closed string sector modes to get a complete description of supersymmetry breaking. Without this there is no goldstino, which signals that the description of supersymmetry breaking is incomplete. It is interesting to ask how the system of D3-branes will evolve once the dilaton and the blow-up mode are allowed to relax and the gravitational back reaction is put in. We leave this question for the future.

Finally, another natural configuration to consider involves not 11, but 8 branes placed at the $\mathbb{Z}_6$ orientifold plane. This is the minimum number required to meet the tadpole conditions (e.g. starting with 32 branes and moving $24 = 6 \times 4$ away leaves us with 8 branes). The corresponding theory has an $SU(4) \times U(1)_A$ gauge symmetry and three generations of fields which transform in the $\begin{smallmatrix}\square\\\square\end{smallmatrix}$ representation of the group. There is in addition an anomalous $U(1)$ under which each of the $\begin{smallmatrix}\square\\\square\end{smallmatrix}$ fields has the same charge. In this case if the FI term (i.e. the orientifold blow-up parameter) vanishes, supersymmetry is unbroken. This is because, in contrast to the $SU(5)$ model, the $SU(4)$ theory has a branch of moduli space where no dynamical superpotential is generated—this can be inferred from [23] by noting that the $SU(4)$ theory with three **6**'s is equivalent to the $SO(6)$ theory with three vectors. The breaking of supersymmetry is then purely due to the $U(1)$ D-term and vanishes for vanishing FI parameter.

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