Brane Supersymmetry Breaking

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Abstract

We show how to construct chiral tachyon-free perturbative orientifold models, where supersymmetry is broken at the string scale on a collection of branes while, to lowest order, the bulk and the other branes are supersymmetric. In higher orders, supersymmetry breaking is mediated to the remaining sectors, but is suppressed by the size of the transverse space or by the distance from the brane where supersymmetry breaking primarily occurred. This setting is of interest for orbifold models with discrete torsion, and is of direct relevance for low-scale string models. It can guarantee the stability of the gauge hierarchy against gravitational radiative corrections, allowing an almost exact supergravity a millimeter away from a non-supersymmetric world.

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The breaking of supersymmetry in String Theory is a long-standing fundamental problem with many ramifications. It is related to the selection of the correct vacuum state, to the cosmological constant problem, to the lifting of flat directions for string moduli, and is a necessary ingredient of a realistic string phenomenology. Unfortunately, despite the recent progress in the understanding of non-perturbative phenomena based on string dualities, little was done on the problem of supersymmetry breaking.

In perturbation theory, closed string vacua with spontaneously broken supersymmetry can be constructed generalizing the Scherk-Schwarz mechanism, and in particular resorting to freely acting orbifolds [1, 2]. The breaking scale is then fixed by a compactification radius, and realistic scenarios ask for radii of the order of a few TeV [2]. This approach is therefore likely to be relevant if the string scale is far below the Planck mass [3], and possibly close to electro-weak energies [4, 5, 6]. A natural framework for such models is the type-I string theory, where gauge interactions are localized on D-branes while gravity propagates in the bulk [5].

Scherk-Schwarz compactifications were recently extended to type I string models in [7], where a new feature was pointed out: the massless spectra of D-branes orthogonal to the coordinate used for the breaking remain supersymmetric at the tree level. As a result, in this case the scale of supersymmetry breaking in the observable world is not directly proportional to the compactification scale, and can have lower values, making this class of constructions more flexible and potentially relevant even if the string scale is moved to intermediate values [8, 9].

The main problem with this mechanism is the cosmological constant. The reason is that the bulk energy density behaves generically as \( \rho_{\text{bulk}} \sim 1/R^{4+d_\perp} \), where \( R \) is the radius of the coordinate used to break supersymmetry and \( d_\perp \) is the dimensionality of the space transverse to our brane world, assumed large with respect to the string scale. The projection

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\(^1\)In the sense that it can be restored by tuning a continuous parameter.
of this cosmological constant on the brane is enhanced by the volume of the transverse space \( r^{d_\perp} \), and is far above \( \mathcal{O}(\text{TeV}^4) \). In fact, the radius \( R \) of a longitudinal direction can not be far from the (TeV) type I string scale \( M_I \) in a perturbative setting, and as a result the brane energy density acquires a quadratic sensitivity to the four-dimensional Planck mass, \( \rho_{\text{brane}} \sim r^{d_\perp} M_I^{4+d_\perp} \sim M_P^2 M_I^2 \) \cite{10}. On the other hand, if \( R \) is transverse \( (R \sim r \ll \mathcal{O}(\text{TeV}^{-1})) \), one obtains \( \rho_{\text{brane}} \sim 1/r^4 \). In both cases, the energy density on the brane is far above the TeV scale, and this destabilizes the hierarchy that one tries to enforce. One way out is to resort to special models with broken supersymmetry but with a vanishing or exponentially small cosmological constant \cite{11}.

Alternatively, one could conceive a different scenario, with supersymmetry broken primordially on our brane world and a string scale of a few TeV. In this case the brane cosmological constant would be, by construction, \( \mathcal{O}(M_I^2) \), while the bulk, only affected by gravitationally suppressed radiative corrections, would be almost supersymmetric \cite{7}. In particular, one would expect that the gravitino mass and the other soft masses in the bulk be extremely small \( \mathcal{O}(M_I^2/M_P) \sim 10^{-4} \text{ eV} \) for \( M_I \sim 1 \text{ TeV} \). Such small masses for scalar moduli and gauge fields might also induce deviations from Newtonian gravity in the (sub)millimeter region that can be experimentally tested \cite{12, 5}. Moreover, the cosmological constant induced in the bulk would be \( \rho_{\text{bulk}} \sim M_I^4/r^{d_\perp} \sim M_I^{6+d_\perp}/M_P^2 \), i.e. of order \( (10 \text{ MeV})^6 \) for \( d_\perp = 2 \). Alternatively, brane supersymmetry breaking could also be of interest in models with an intermediate string scale \( \sim 10^{11} \text{ GeV} \) \cite{9}, if it occurs on a brane distant from our world and is therefore mediated to us by gravitational interactions.

The purpose of this letter is to show that it is possible to construct perturbative orientifold models \cite{13} where supersymmetry breaking originates from a collection of branes, while both the bulk spectrum and the spectrum of other branes are supersymmetric at tree-level. Whereas models with bulk supersymmetry can naturally be constructed in the effective field theory, for instance appealing to non-perturbative super-Yang-Mills dynam-
ics, we believe that a direct string construction is of some interest, in particular to attain a better comparison with field theory supersymmetry breaking mechanisms.

Brane supersymmetry breaking can also be induced turning on internal magnetic fields \[14\]. The mechanism we are proposing shares some properties with this setting, although supersymmetry is broken at the string scale, but appears to avoid some of its problems, namely the presence of tachyons and the generic lack of gaugino masses for the unbroken gauge group.

The rules for constructing perturbative type-I orientifolds \[15, 16\] rest on the modular invariance of the closed string spectrum and on some conditions linking its Klein-bottle projection to the open and unoriented sector. These are to be supplemented by Ramond-Ramond (RR) tadpole conditions, that are directly related to anomaly cancellations and may be regarded as global neutrality conditions for RR charges in a compact internal space \[17\]. Tadpole cancellations result from opposite contributions of boundaries and crosscaps or, equivalently, of branes and orientifolds.

It is possible to construct models where some of the orientifold contributions are inverted, so that the necessary RR cancellations require some care. A simple example of this phenomenon is afforded by the \(T^4/Z_2\) orientifold where, compatibly with the “crosscap constraint” \[18\], all twisted Klein-bottle contributions are reversed. These exotic Klein-bottle projections are quite interesting, and have already led to tachyon-free non-supersymmetric open-string vacua \[18\]. In our case, the resulting unoriented closed spectrum is supersymmetric and contains, aside from the \((1,0)\) gravitational multiplet, 17 tensor multiplets and 4 hypermultiplets, but its RR tadpoles can not be canceled in the usual way. A more sophisticated, but physically very interesting, set of examples, is provided by the open descendants of \(Z_2 \times Z_2\) models, and in particular of those with discrete torsion, where some of the orientifold charges are necessarily reversed. In all these cases, the tadpole conditions may be solved unpairing NS and R contributions, and thus inducing brane supersymmetry
breaking. In this letter, for the sake of brevity, we confine our attention to the $T^4/Z_2$ case, where one obtains a chiral 6D spectrum that is free of tachyons and satisfies all usual anomaly cancellation constraints. In this model, that contains 32 D9 and 32 anti-D5 branes, the absence of tachyons can also be understood as in recent studies \cite{20} of stable non-BPS states of the type IIB string.

The same mechanism can be applied to the $Z_2 \times Z_2$ models, that will be discussed elsewhere, and in principle should offer the possibility to deal with other orientifold models where tadpole conditions do not admit naive supersymmetric solutions.

### A six-dimensional example

We now present an explicit 6D model containing D9 and D5 (anti-D5) branes, where at tree level supersymmetry breaking is induced on the 55 and 59 states, while the 99 states and the bulk (closed) spectrum are supersymmetric. The starting point in this construction is a modification of the $\Omega$ projection in the twisted sector of the $T^4/Z_2$ model. This inverts the charge of the O5 planes, and is actually compatible with the perturbative rules of orientifold models, but the cancellation of RR tadpoles requires 32 D5 branes. Omitting for brevity the contributions of the transverse bosons, the torus partition function is

$$
\mathcal{Z} = \frac{1}{2} |Q_o + Q_v|^2 \Lambda + \frac{1}{2} |Q_o - Q_v|^2 \frac{2\eta^4}{\theta^2} + \frac{1}{2} |Q_s + Q_c|^2 \frac{2\eta^4}{\theta_4^2} + \frac{1}{2} |Q_s - Q_c|^2 \frac{2\eta^4}{\theta_3^2},
$$

where $\Lambda$ denotes the compactification lattice. In $\mathcal{Z}$, we have introduced the convenient combinations of $SO(4)$ characters

$$
Q_o = V_4 O_4 - C_4 C_4, \quad Q_v = O_4 V_4 - S_4 S_4,
$$

$$
Q_s = O_4 C_4 - S_4 O_4, \quad Q_c = V_4 S_4 - C_4 V_4,
$$

defined as

$$
O_4 = \frac{\theta_3^2 + \theta_4^2}{2\eta^2}, \quad V_4 = \frac{\theta_3^2 - \theta_4^2}{2\eta^2}, \quad S_4 = \frac{\theta_2^2 - \theta_1^2}{2\eta^2}, \quad C_4 = \frac{\theta_2^2 + \theta_1^2}{2\eta^2}.
$$
Turning to the Klein bottle, let us consider the two inequivalent choices
\[ \mathcal{K} = \frac{1}{4} \left\{ (Q_o + Q_v)(P + W) + 2\epsilon \times 16(Q_s + Q_c)\left( \frac{\eta}{\theta_4} \right)^2 \right\}, \]
\[ (3) \]
where \( P \) (\( W \)) denotes the momentum (winding) lattice sum and \( \epsilon = \pm 1 \). For both choices of \( \epsilon \), the closed string spectrum has \( (1, 0) \) supersymmetry, but the two resulting projections are quite different. The usual choice \( (\epsilon = 1) \) leaves 1 gravitational multiplet, 1 tensor multiplet and 20 hypermultiplets, while \( \epsilon = -1 \) leaves 1 gravitational multiplet, 17 tensor multiplets and 4 hypermultiplets\(^3\). The projected closed spectrum coincides with the one considered in ref. [21], but our open spectrum is markedly different. The transverse-channel Klein bottle amplitude reads
\[ \tilde{\mathcal{K}} = \frac{2^5}{4} \left\{ (Q_o + Q_v) \left( vW + \frac{P^e}{v} \right) + 2\epsilon (Q_o - Q_v) \left( \frac{2\eta}{\theta_2} \right)^2 \right\}, \]
\[ (4) \]
where \( P^e \) (\( W^e \)) denotes the lattice of even momenta (windings) and \( v \) is the volume of the compact space. The reversal of the O5 charge respects the positivity structure at the origin of the \( T^4 \) lattice, and indeed the coefficients combine into perfect squares:
\[ \tilde{\mathcal{K}}_0 = \frac{2^5}{4} \left\{ Q_o \left( \sqrt{v} + \frac{\epsilon}{\sqrt{v}} \right)^2 + Q_v \left( \sqrt{v} - \frac{\epsilon}{\sqrt{v}} \right)^2 \right\}. \]
\[ (5) \]
The case \( \epsilon = 1 \) leads to the familiar \( U(16) \times U(16) \) model [13, 16], and therefore from now on we restrict our attention to the other choice.

RR tadpole cancellations and the modifications of \( \mathcal{K} \) are compatible with the positivity of the open spectrum in \( \mathcal{A} \), provided one introduces D5 branes, rather than the usual D5 branes. In the transverse channel, this choice inverts the signs of all RR contributions in the Neumann-Dirichlet (95) part, but leaves all other untwisted terms unchanged. Introducing suitable Chan-Paton charges \( N, D \) and their \( Z_2 \) orbifold breakings \( R_N, R_D \), the transverse annulus amplitude reads
\[ \tilde{\mathcal{A}} = \frac{2^{-5}}{4} \left\{ (Q_o + Q_v) \left( N^2vW + \frac{D^2P}{v} \right) + 2ND(Q_o' - Q_v') \left( \frac{2\eta}{\theta_2} \right)^2 \right\}. \]
\[ (6) \]
\(^2\)The world-sheet moduli for the various amplitudes, implicit in the following, are defined as in [7].
\(^3\)A similar projection with a non-zero \( B_{ab} \) [15] would result in 13 or 11 tensor multiplets [22].
+ 16(Q_s + Q_c) \left( R_N^2 + R_D^2 \right) \left( \frac{\eta}{\theta_2} \right)^2 + 8 R_N R_D (V_4 S_4 - O_4 C_4 - S_4 O_4 + C_4 V_4) \left( \frac{\eta}{\theta_3} \right)^2 \right),

where we have also introduced primed characters, related by a chirality change \( S_4 \leftrightarrow C_4 \) to the unprimed ones in \((\mathbb{I})\):

\[
\begin{align*}
Q'_o &= V_4 O_4 - S_4 S_4, \\
Q'_v &= O_4 V_4 - C_4 C_4, \\
Q'_s &= O_4 S_4 - C_4 O_4, \\
Q'_c &= V_4 C_4 - S_4 V_4.
\end{align*}
\tag{7}
\]

As usual, at the origin of the lattice the different terms organize into perfect squares:

\[
\bar{A}_0 = \frac{2^{-5}}{4} \left\{ Q'_o \left( N \sqrt{\bar{v}} + \frac{D}{\sqrt{\bar{v}}} \right)^2 + Q'_v \left( N \sqrt{\bar{v}} - \frac{D}{\sqrt{\bar{v}}} \right)^2 \right. \\
+ \left. (V_4 S_4 - S_4 O_4) \left( 15 R_N^2 + (R_N + 4 R_D)^2 \right) + (O_4 C_4 - C_4 V_4) \left( 15 R_N^2 + (R_N - 4 R_D)^2 \right) \right\}.
\tag{8}
\]

Here, for simplicity, all D5 branes, whose geometry is neatly displayed by the breaking terms, have been placed at the origin of the compact space. The direct-channel annulus is obtained by an S-transformation, and reads

\[
\bar{A} = \frac{1}{4} \left\{ (Q_o + Q_v)(N^2 P + D^2 W) + 2 N D (Q'_o + Q'_v) \left( \frac{\eta}{\theta_1} \right)^2 \right. \\
+ \left. (R_N^2 + R_D^2)(Q_o - Q_v) \left( \frac{2 \eta}{\theta_2} \right)^2 + 2 R_N R_D (-O_4 S_4 - C_4 O_4 + V_4 C_4 + S_4 V_4) \left( \frac{\eta}{\theta_3} \right)^2 \right\}.
\tag{9}
\]

Finally, the Möbius amplitude at the origin of the lattices (for the definition of the hatted characters, see \([\mathbb{II}]\)),

\[
\tilde{M}_0 = -\frac{1}{2} \left[ \hat{V}_4 \hat{O}_4 \left( \sqrt{\bar{v}} - \frac{1}{\sqrt{\bar{v}}} \right) \left( N \sqrt{\bar{v}} + \frac{D}{\sqrt{\bar{v}}} \right) + \hat{O}_4 \hat{V}_4 \left( \sqrt{\bar{v}} + \frac{1}{\sqrt{\bar{v}}} \right) \left( N \sqrt{\bar{v}} - \frac{D}{\sqrt{\bar{v}}} \right) \right. \\
- \left. \hat{C}_4 \hat{C}_4 \left( \sqrt{\bar{v}} - \frac{1}{\sqrt{\bar{v}}} \right) \left( N \sqrt{\bar{v}} + \frac{D}{\sqrt{\bar{v}}} \right) + \hat{S}_4 \hat{S}_4 \left( \sqrt{\bar{v}} + \frac{1}{\sqrt{\bar{v}}} \right) \left( N \sqrt{\bar{v}} + \frac{D}{\sqrt{\bar{v}}} \right) \right],
\tag{10}
\]

is easily obtained combining \( \tilde{K}_0 \) and \( \bar{A}_0 \), and allows one to reconstruct the full transverse Möbius amplitude

\[
\tilde{M} = -\frac{1}{2} \left\{ N v W^e (\hat{V}_4 \hat{O}_4 + \hat{O}_4 \hat{V}_4 - \hat{C}_4 \hat{C}_4 - \hat{S}_4 \hat{S}_4) - \frac{D P^e}{v} (\hat{V}_4 \hat{O}_4 + \hat{O}_4 \hat{V}_4 + \hat{C}_4 \hat{C}_4 + \hat{S}_4 \hat{S}_4) \right. \\
- N (\hat{V}_4 \hat{O}_4 - \hat{O}_4 \hat{V}_4 - \hat{C}_4 \hat{C}_4 + \hat{S}_4 \hat{S}_4) \left( \frac{2 \eta}{\theta_2} \right)^2 + D (\hat{V}_4 \hat{O}_4 - \hat{O}_4 \hat{V}_4 + \hat{C}_4 \hat{C}_4 - \hat{S}_4 \hat{S}_4) \left( \frac{2 \eta}{\theta_2} \right)^2 \right\}.
\tag{11}
\]
Eq. (11) shows some marked differences with respect to the more familiar model with $\epsilon = 1$. These may be given a neat physical interpretation, since $\tilde{M}$ describes the propagation between holes and crosscaps, or equivalently between branes and orientifold planes. Therefore, one can see that all D9-O9 terms, the D5-O5 terms in the R-R sector and the D5-O9 terms in the NS-NS sector are as in the standard $T^4/Z_2$ orientifold, while the signs of all D9-O5 terms, of the D5-O5 terms in the NS-NS sector and of the D5-O9 terms in the R-R sector are inverted. In particular, this implies that the M"obius amplitude breaks supersymmetry at tree level in the D5 sector, an effect felt by all open-strings ending on the D5 branes.

Finally, a P transformation determines the direct (open string) amplitude

$$\mathcal{M} = -\frac{1}{4}\left\{NP(\hat{O}_4 \hat{V}_4 + \hat{V}_4 \hat{O}_4 - \hat{S}_4 \hat{S}_4 - \hat{C}_4 \hat{C}_4) - DW(\hat{O}_4 \hat{V}_4 + \hat{V}_4 \hat{O}_4 + \hat{S}_4 \hat{S}_4 + \hat{C}_4 \hat{C}_4)
- N(\hat{O}_4 \hat{V}_4 - \hat{V}_4 \hat{O}_4 - \hat{S}_4 \hat{S}_4 + \hat{C}_4 \hat{C}_4) \left(\frac{2\bar{\eta}}{\theta_2}\right)^2 + D(\hat{O}_4 \hat{V}_4 - \hat{V}_4 \hat{O}_4 + \hat{S}_4 \hat{S}_4 - \hat{C}_4 \hat{C}_4) \left(\frac{2\bar{\eta}}{\theta_2}\right)^2 \right\} \ . (12)$$

Parametrizing the Chan-Paton charges as

$$N = n_1 + n_2 \ , \quad D = d_1 + d_2 \ , \quad R_N = n_1 - n_2 \ , \quad R_D = d_1 - d_2 \ , \quad (13)$$

the RR tadpole conditions $N = D = 32, R_N = R_D = 0 \ (n_1 = n_2 = d_1 = d_2 = 16)$ determine the gauge group $[SO(16) \times SO(16)]_9 \times [USp(16) \times USp(16)]_5$.

The massless matter representations may be read from:

$$\mathcal{A}_0 + \mathcal{M}_0 = \frac{n_1(n_1 - 1) + n_2(n_2 - 1) + d_1(d_1 + 1) + d_2(d_2 + 1)}{2} V_4 O_4
- \frac{n_1(n_1 - 1) + n_2(n_2 - 1) + d_1(d_1 - 1) + d_2(d_2 - 1)}{2} C_4 C_4
+ (n_1n_2 + d_1d_2)(O_4 V_4 - S_4 S_4) + (n_1d_2 + n_2d_1) O_4 S_4 - (n_1d_1 + n_2d_2) C_4 O_4 \ . \quad (14)$$

The 99 spectrum is supersymmetric, and comprises the (1,0) vector multiplets for the $SO(16) \times SO(16)$ gauge group and a hypermultiplet in the representations $(16, 16, 1, 1)$ of the gauge group. On the other hand, the 55 DD spectrum is not supersymmetric, and
contains, aside from the gauge vectors of $[USp(16) \times USp(16)]$, quartets of scalars in the $(1, 1, 16, 16)$, right-handed Weyl fermions in the $(1, 1, 120, 1)$ and in the $(1, 1, 1, 120)$, and left-handed Weyl fermions in the $(1, 1, 16, 16)$. Finally, the ND sector is also non supersymmetric, and comprises doublets of scalars in the $(16, 1, 16)$ and in the $(1, 16, 16, 1)$, together with additional (symplectic) Majorana-Weyl fermions in the $(16, 1, 16, 1)$ and $(1, 16, 1, 16)$. These Majorana-Weyl fermions are a peculiar feature of six-dimensional space time, where the fundamental Weyl fermion, a pseudoreal spinor of $SU^*(4)$, can be subjected to an additional Majorana condition, if this is supplemented by a conjugation in a pseudoreal representation $[23]$. In this case, this is indeed possible, since the ND fermions are valued in the fundamental representation of $USp(16)$. This doubling is also a useful technical trick in six-dimensional supergravity, where Weyl fermions are often extended to $Sp(2)$ Majorana-Weyl doublets in the $(1, 0)$ case and to $Sp(4)$ Majorana-Weyl quartets in the $(2, 0)$ case.

It should be appreciated that, from the $D9$ brane point of view, the diagonal combination of the two $USp(16)_5$ gauge groups acts as a global symmetry. This corresponds to having complex scalars and symplectic Majorana-Weyl fermions in the representations $16 \times [(16, 1) + (1, 16)]$ of the D9 gauge group. As a result, the non-supersymmetric ND spectrum looks effectively supersymmetric, and indeed all 95 terms do not contribute to the vacuum energy. However, as in the 6D temperature breaking discussed in [7], the chirality of the fermions in $Q'_s$ is not the one required by 6D supersymmetry. This chirality flip is a peculiar feature of six-dimensional models, and does not persist in the reduction to four dimensions [1]. Thus, in four dimensions supersymmetry would be strictly unbroken on the $D9$ branes before turning on $D5$ gauge interactions. This setting shows some marked

\footnote{In view of the present results, the 6D model presented in the first reference in [7] actually contains, in addition to the $D9$ branes, one set of 16 $D5$ branes and one set of 16 $D\bar{5}$ branes. This justifies both the chirality changes in the twisted spectrum of the $D5$ branes and the presence of the tachyonic mode in the 5$\bar{5}$ annulus amplitude.}
differences with respect to the mechanism of bulk supersymmetry breaking (or brane supersymmetry), where the massless ND sector was supersymmetric from both the N and D viewpoints [7].

It is easy to check that all the irreducible gauge and gravitational anomalies cancel in this model as a result of the tadpole conditions for the RR fields. The residual anomaly polynomial does not factorize, and reveals the need for a generalized Green-Schwarz mechanism [24], with couplings of a more general type than those found in supersymmetric models:

\[
\begin{align*}
A &= \frac{1}{64}(F_1^2 + F_2^2 + F_3^2 + F_4^2 - 2R^2)^2 + \frac{3}{64}(F_1^2 - F_2^2 - F_3^2 + F_4^2)^2 \\
&+ \frac{5}{64}(F_1^2 - F_2^2 + F_3^2 - F_4^2)^2 - \frac{1}{64}(F_1^2 + F_2^2 - F_3^2 - F_4^2)^2 .
\end{align*}
\]

(15)

Brane-antibrane interactions have been discussed recently in the literature in the context of stable non-BPS states [20]. Our results for the D9-D5 system, restricted to the open oriented sector, provide particular examples of type-I vacua including stable non-BPS states with vanishing interaction energy for all radii, as can be seen from the vanishing of the ND annulus amplitude.

Comparing the unoriented closed and open string amplitudes (3), (9) and (12) with those of the supersymmetric \( T^4/Z_2 \) orientifold, it is easy to see that the former can be obtained from the latter by a \( \pi \)-rotation of the D5 branes together with a reversal of the charge of the O5 planes. This phenomenon presents some analogies with the deformations induced by an internal magnetic field [25] felt only by the D5 branes. This would shift the oscillator mode numbers of the 55 strings by an amount \( \epsilon = \epsilon_I + \epsilon_J \), where

\[
\epsilon_I = \frac{1}{\pi} \arctan(\pi q_I H) ,
\]

(16)

with \( q_I \) the charges of the corresponding Chan-Paton states. In a similar fashion, it would shift the mode numbers of the 95 strings by \( \epsilon = \epsilon_I \), and the Möbius contributions by \( \epsilon = 2\epsilon_I \). These oscillator shifts translate into the mass shifts \( \delta M^2 = (2n + 1)|\epsilon| + 2\epsilon S_{int} \),
where the integer \( n \) denotes the Landau level and \( S_{\text{int}} \) denotes the internal spin. As a result, in a generic magnetic field the internal components of the higher-dimensional gauge fields become tachyonic. This is not the only problem presented by type-I strings in magnetic fields: RR tadpoles are generally non-vanishing while, from a more phenomenological perspective, the masses of all the gauginos of the unbroken gauge group, that are neutral under \( Q \), vanish.

Our mechanism consists partly of introducing a discrete value \( \epsilon_I = 1 \), outside the physical range of eq. (16). As a result, in the annulus the 99 and 55 contributions are unaffected, while in the 95 terms the internal characters are interchanged, according to \( S_4 \leftrightarrow C_4 \) and \( O_4 \leftrightarrow V_4 \). On the other hand, as we have seen, the modifications of the Möbius amplitude are more subtle, due to the simultaneous sign change of the O5 charge, that results into a symplectic gauge group in the 55 sector with no massless adjoint fermions. Thus, the spectrum is tachyon free, all RR tadpoles are canceled and there are no massless gauginos.

The breaking of supersymmetry gives rise to a vacuum energy localized on the \( D_5 \) branes, and thus to a tree-level potential for the NS moduli, that can be extracted from the corresponding uncanceled NS tadpoles in eqs. (3), (8) and (10). A simple inspection shows that the only non-vanishing ones correspond to the NS characters \( V_4O_4 \) and \( O_4V_4 \) associated to the 6D dilaton \( \phi_6 \) and to the internal volume \( v \):

\[
\frac{2^{-5}}{4} \left\{ \left( (N - 32) \sqrt{v} + \frac{D + 32}{\sqrt{v}} \right)^2 V_4O_4 + \left( (N - 32) \sqrt{v} - \frac{D + 32}{\sqrt{v}} \right)^2 O_4V_4 \right\} .
\]

Using factorization and the values \( N = D = 32 \) needed to cancel the RR tadpoles, the potential (in the string frame) is:

\[
V_{\text{eff}} = c \frac{e^{-\phi_6}}{\sqrt{v}} = ce^{-\phi_{10}} = \frac{c}{g_{\text{YM}}^2},
\]

where \( \phi_{10} \) is the 10D dilaton, that determines the Yang-Mills coupling \( g_{\text{YM}} \) on the D5 branes, and \( c \) is some positive numerical constant. The potential (18) is clearly localized
on the D5 branes, and is positive. This can be understood noticing that the O9 plane contribution to vacuum energy, identified from (17), is negative and exactly cancels for $N = 32$. This fixes the D5 brane contribution to the vacuum energy, that is thus positive, consistently with the interpretation of this mechanism as global supersymmetry breaking. The potential (18) has the usual runaway behavior, as expected by general arguments.

An interesting application of the above mechanism is to the $Z_2 \times Z_2$ orbifold model with discrete torsion, where nontrivial two-form fluxes at the orbifold fixed points invert the sign of the lattice independent terms in the torus amplitude. It is interesting to notice that the corresponding supersymmetric spectra, which contain chiral fermions, are inconsistent because of un-cancelled RR tadpoles. The solution to the RR tadpoles can again be obtained introducing anti-D5 branes, along the lines described previously. We have worked out in detail the partition functions and the spectra of the open descendants, that will be presented elsewhere. Here we conclude by describing some qualitative features of the resulting models. The O5$_i$ plane charges, manifest in the Klein bottle amplitude, are equal to $-32\epsilon_i$, where the signs $\epsilon_i = \pm 1$ are restricted by $\epsilon_1\epsilon_2\epsilon_3 = -1$. There are thus four independent possibilities, $(\epsilon_1, \epsilon_2, \epsilon_3) = (-1, -1, -1), (-1, 1, 1), (1, -1, 1)$ and $(1, 1, -1)$. As in the 6D model discussed in the present paper, every O5$_i$ charge flip $\epsilon_i = -1$ will ask for a set of 32 D¯5$_i$ branes. Moreover, in this case the D5$_i$ gauge group becomes unitary, whereas for $\epsilon_i = 1$ it is symplectic as in the supersymmetric model without discrete torsion. However, the gauginos of unitary groups are massless, and can acquire masses only through quantum corrections. Indeed, in our 6D example the gauginos were massive because the Möbius amplitude had different projections for bosons and fermions. On the other hand, the adjoint representation for unitary groups is not affected by the Möbius amplitude, and therefore the corresponding gauginos stay massless. The D9 gauge group changes too: it is $SO(8)\mathbb{C}_4^2$ for $(\epsilon_1, \epsilon_2, \epsilon_3) = (-1, -1, -1)$ and $U(8)\mathbb{C}_2^2$ for the other 3 choices. Supersymmetry is broken at the string scale in the sectors $9\bar{5}_i$ and in the Dirichlet part of the Möbius amplitude, in analogy with the 6D model described here. Moreover, if $\epsilon_i = 1$ and $\epsilon_j = -1$
supersymmetry is also broken in the sector $5_i \bar{5}_j$, where a tachyon appears, in agreement with the analysis of non-BPS states of [20]. Finally, the remaining NS-NS tadpoles give rise to a scalar potential localized on the anti-D5 branes, as in eq. (18).

It would be interesting to investigate along the same lines other models, as the 4d $\mathbb{Z}_4$ orientifold, where the tadpole conditions have no naive solution.

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