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# Tachyon-free non-supersymmetric type IIB orientifolds via brane-antibrane systems

**Gerardo Aldazabal**

*Instituto Balseiro, CNEA, Centro Atómico Bariloche,*
*8400 S.C. de Bariloche, and CONICET, Argentina.*
*E-mail:* aldazaba@ictp.trieste.it

**Angel M. Uranga**

*School of Natural Sciences, Institute for Advanced Study*
*Olden Lane, Princeton NJ 08540, USA*
*E-mail:* uranga@IAS.EDU

Abstract: We derive the rules to construct type IIB compact orientifolds in six and four dimensions including D-branes and anti-D-branes. Even though the models are non-supersymmetric due to the presence of the anti-D-branes, we show that it is easy to construct large classes of models free of tachyons. Brane-antibrane annihilation can be prevented for instance by considering models with branes and antibranes stuck at different fixed points in the compact space. We construct several anomaly-free and tachyon-free six-dimensional orientifolds containing D9-branes and anti-D5-branes. This setup allows to construct four-dimensional chiral models with supersymmetry unbroken in the bulk and in some D-brane sectors, whereas supersymmetry is broken (at the string scale) in some 'hidden' anti-D-brane sector. We present several explicit models of this kind. We also comment on the role of the non-cancelled attractive brane-antibrane forces and the non-vanishing cosmological constant, as providing interesting dynamics for the geometric moduli and the dilaton, which may contribute to their stabilization.

Keywords: D-branes, Superstring Vacua, Supersymmetry Breaking.

## Contents



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# 1. Introduction

The last two years have witnessed an increased interest on non-supersymmetric string theory compactifications and non-supersymmetric states in string theory. This has led to a number of interesting developments. For instance, the construction of non-supersymmetric string vacua with vanishing or extremely suppressed cosmological constant [1]. Another example is the study of tachyon condensation in brane-antibrane systems and its relation to stable non-BPS states in string theory [2]. Hopefully these and other developments will provide new insights into the possible mechanisms breaking supersymmetry in realistic compactifications of string theory.

One of the most interesting classes of string theory vacua is that of type IIB orientifolds [3]–[9]. Supersymmetric compact orientifolds in six [6]–[12] and four [13, 14, 15, 16, 9] dimensions have been constructed and studied extensively. They show several interesting properties, both from the theoretical and the phenomenological points of view. Within the latter, we may highlight the fact that they provide explicit string theory models where gauge fields and charged matter are localized on D-branes, whereas gravity and other closed string modes propagate in full spacetime. This idea has recently received a great deal of attention, in particular in the context of compactifications with large [17, 18], intermediate [19] or 'largish' [20] compact dimensions. It is natural to wonder about possible supersymmetry breaking mechanisms in these scenarios.

Several possibilities to achieve the construction of non-supersymmetric orientifold vacua have been explored in the literature. For instance, one possibility is to implement the Scherk-Schwarz mechanism [21] in the context of type IIB orientifolds [22]. A different approach would consist of constructing orientifolds of non-supersymmetric string theories, as in type 0B orientifolds in [23].

In this paper we would like to point out a different possibility. Non-supersymmetric type IIB orientifolds can be easily obtained by introducing anti-D-branes in the model. This framework has several built-in advantages. For example, since the orientifold projections on the closed string spectrum preserve some of the supersymmetries, closed string tachyons are automatically absent. Perhaps a bit more surprisingly, open string tachyons (arising from open strings stretched between branes and antibranes) are also easily avoided, for instance by choosing Chan-Paton embeddings that project out the tachyon of coincident branes and antibranes, or by locating the branes and antibranes stuck at different fixed points in the internal space. Therefore, large classes of tachyon-free non-supersymmetric type IIB orientifolds can be constructed in six and four dimensions. A nice feature that arises in some four-dimensional models is that they can include chiral $\mathcal{N} = 1$ supersymmetric sectors of D-branes and non-supersymmetric sectors of anti-D-branes, spatially separated in the compact space. This leads to explicit string theory models where supersymmetry breaking (at the string scale) occurs in a distant brane and is transmitted to the visible supersymmetric sector, either by bulk modes (gravity-mediated models) or by a gauge sector on other D-branes coupling to both branes and antibranes (gauge-mediated).

These models also present interesting properties from the theoretical point of view. For instance we show that certain non-supersymmetric compactification of type I on orbifold limits of K3 are related by T-duality to analogous compactification of the ten-dimensional non-supersymmetric USp(32) theory constructed in [24].

The paper is organized as follows. In section 2 we provide the basic rules to construct type IIB orbifold and orientifold models including antibranes, and compute the spectrum of tachyons and massless states. This allows to derive conditions under which the open string tachyons are absent. We also discuss the RR tadpole cancellation conditions for these models.

In section 3 we present several examples of non-supersymmetric type IIB orientifolds in six dimensions. We discuss how some naively stable models, without tachyons but in which the antibranes can move off the fixed points into the bulk, can actually decay to supersymmetric theories through brane-antibrane annihilation. We also show how the introduction of antibranes in some cases allows to satisfy tadpole conditions which cannot be fulfilled in supersymmetric models with only D-branes. We illustrate this property by constructing orientifolds of type IIB on $T^4/\mathbb{Z}_N$, $N = 2, 4, 6$, with vector structure.

In section 4, we present four-dimensional examples based on the $\mathbb{Z}_6'$ and $\mathbb{Z}_3$ orientifolds, which have unbroken $\mathcal{N} = 1$ supersymmetry in the closed string sector and in several D-brane sectors, and non-supersymmetric sectors of anti-D-branes.

Section 5 contains our final remarks and speculations concerning the role of the non-supersymmetric dynamics in stabilizing the geometric moduli and the dilaton in these models.

While we were completing the present paper, we noticed reference [25], where one of the models presented in section 3.1 is constructed using slightly different techniques.

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## 2. General construction

### 2.1 Generalities of type IIB orientifolds

In this section we summarize the basic ingredients [3, 6, 7, 9] and notation needed for the construction of $N = 1$ type IIB orientifolds.[1]

A type IIB orbifold is obtained when the toroidally compactified theory is divided out by discrete symmetry group $G_1$, (like $\mathbb{Z}_N$ or $\mathbb{Z}_N \times \mathbb{Z}_M$). The orbifold twist eigenvalues $v_a$, associated to a complex compact coordinate $Y_a$ $(a = 1, \ldots, D - 10)$, are restricted by the requirement that the orbifold group acts crystalographically on $T^{D-10}$ and by the number of supersymmetries to be left unbroken. For instance for $\mathbb{Z}_N$ in $D = 6$ and $\mathcal{N} = 2$ unbroken supersymmetries we must have $v = \frac{1}{N}(1, -1)$. The same requirement in $D = 4$ is satisfied by $v = \frac{1}{N}(\ell_1, \ell_2, \ell_3)$ with $\ell_1 + \ell_2 + \ell_3 = 0$, where $\ell_a$ are some specific integers [26].

A type IIB orientifold results from the joint action of the orbifold group $G_1$ together with a world sheet parity operation $\Omega$, exchanging left and right movers. By keeping orientation reversal invariant states, supersymmetry is reduced by a half. $\Omega$ action can be also accompanied by extra operations thus leading to generic orientifold group $G_1 + \Omega G_2$ with $\Omega h \Omega h' \in G_1$ for $h, h' \in G_2$

Orientifolding closed type IIB string introduces a Klein-bottle unoriented world-sheet. Amplitudes on such a surface contain tadpole divergences. Tadpoles may be

[1] In the orientifolds we consider in coming sections, supersymmetry is preserved in the closed string sector, and is only broken in the open string sector by the presence of the anti-D-branes. Hence most of the usual techniques in supersymmetric orientifolds apply.

generically interpreted as unbalanced orientifold plane charges under RR form potentials. In order to eliminate such unphysical divergences Dp-branes, which carry opposite charges, must be generically introduced. In this way, divergences occurring in the open string sector cancel up the closed sector ones and produce a consistent theory.

We will focus our discussion in $D = 6, 4$ dimensional theories.

For $\mathbb{Z}_N$, with $N$ odd, only D9-branes are required. They fill the full space-time and the compact space. For $N$ even, D5-branes are required. Their world-volume fills up the space-time in six dimensions.

D$5_k$-branes, with world volume filling space-time and the $k^{th}$ complex plane, may be required. This is so whenever the orientifold group contains the element $\Omega R_i R_j$, for $k \neq i, j$. Here $R_i$ ($R_j$) is an order two twist of the $i^{th}$ ($j^{th}$) complex plane.

Open string states are denoted by $|\Psi, ab\rangle$, where $\Psi$ refers to world-sheet degrees of freedom while the $a, b$ Chan-Paton indices are associated to the open string endpoints lying on D$p$-branes and D$q$-branes respectively.

These Chan-Paton labels must be contracted with a hermitian matrix $\lambda^{pq}_{ab}$. The action of an element of the orientifold group on Chan-Paton factors is achieved by a unitary matrix $\gamma_{g,p}$ such that $g : \lambda^{pg} \to \gamma_{g,p}\lambda^{pq}\gamma^{-1}_{g,q}$. We denote by $\gamma_{k,p}$ the matrix associated to the $Z_N$ orbifold twist $\theta^k$ acting on a D$p$-brane.

Consistency under group transformations imposes restrictions on the representations $\gamma_g$. For instance, from $\Omega^2 = 1$ it follows that

$$\gamma_{\Omega,p} = \pm\gamma^T_{\Omega,p} \, . \tag{2.1}$$

Also group operation $\left(\Omega\theta^k\right)^2 = \theta^{2k}$ and (2.3) lead to

$$\gamma^*_{k,p} = \pm\gamma^*_{\Omega,p}\,\gamma_{k,p}\gamma_{\Omega,p} \, , \tag{2.2}$$

for $p = 9, 5$.

Tadpole cancellation imposes further constraints on $\gamma_g$ (see for instance [8, 9] and references therein).

In what follows we will make a definite choice of signs in (2.1), namely

$$\begin{aligned} \gamma_{\Omega,9} &= \gamma^T_{\Omega,9} \\ \gamma_{\Omega,5} &= -\gamma^T_{\Omega,5} \end{aligned} \tag{2.3}$$

for $\Omega$ acting on 9 and on 5-branes. The first condition is the usual requirement of global consistency of the ten-form potential in type I theory. Second equation is in agreement with the Gimon and Polchinski action, analyzed in [7].

Generic matrices satisfying above constraints can be provided. Namely, for a $Z_N$ orbifold twist action, with $N = 2P$ ($N = 2P + 1$) we define

$$\gamma_{1,p} = (\tilde{\gamma}_{1,p}, \tilde{\gamma}^*_{1,p}) \, , \tag{2.4}$$

with $*$ denoting complex conjugation and where $\tilde{\gamma}$ is a $N_p \times N_p$ diagonal matrix given by

$$\tilde{\gamma}_{1,p} = \mathrm{diag}(\ldots, \alpha^{NV_j} I_{n_j^p}, \ldots, \alpha^{NV_P} I_{n_P^p}) \tag{2.5}$$

with $\alpha = \mathrm{e}^{2i\pi/N}$ and $2N_p = 2\sum_{j=1}^{P} n_j^p$ the number of D$p$-branes.

The choice $V_j = j/N$ with $j = 0, \ldots, P$ corresponds to an action "with vector structure " ($\gamma^N = 1$) while $V_j = \frac{2j-1}{2N}$ with $j = 1, \ldots, P$ describes an action "without vector structure' ($\gamma_{1,p}^N = -1$).[2]

By choosing $\gamma_{\Omega,9}$ and $\gamma_{\Omega,5}$ matrices

$$\gamma_{\Omega,9} = \begin{pmatrix} 0 & \mathrm{I}_{N_9} \\ \mathrm{I}_{N_9} & 0 \end{pmatrix}, \quad \gamma_{\Omega,5} = \begin{pmatrix} 0 & -i\mathrm{I}_{N_5} \\ i\mathrm{I}_{N_5} & 0 \end{pmatrix}, \tag{2.6}$$

then (2.3) and (2.2) (with $+(-)$ sign for D9(5)-brane respectively) are satisfied.

The open string spectrum is constructed by requiring the states $|\Psi, ab\rangle\lambda_{ab}$ to be invariant under the action of the orientifold group. We present the general rules in the following section.

### 2.2 Open string spectrum with branes and antibranes

In this section we discuss the rules to construct the spectrum of tachyonic and massless states in the open string sector, in orientifolds containing both D-branes and anti-D-branes (denoted $\overline{\mathrm{D}}$-branes in what follows). Even though several of the models we are going to construct are six-dimensional, it will be convenient to derive the spectrum for a general four-dimensional twist $v = (v_1, v_2, v_3) = \frac{1}{N}(\ell_1, \ell_2, \ell_3)$. The rules for six-dimensional models are recovered by setting $v_3 = 0$ and $v_1 = -v_2 = 1/N$, and by taking into account that there are two additional spacetime dimensions.

We start by reviewing the supersymmetric sectors.

**The 99 sector** In the NS sector the GSO projection eliminates the tachyon. Massless gauge bosons and complex scalars in spacetime are obtained from the states $\psi^\mu_{-\frac{1}{2}}|0, ab\rangle\lambda^{(0)}_{ab}$, $\psi^i_{-\frac{1}{2}}|0, ab\rangle\lambda^{(i)}_{ab}$, where $i = 1, 2, 3$ labels the three complex planes in the orientifold. The projections on the Chan-Paton factors of these states are

$$\begin{aligned} \lambda^{(0)} &= \gamma_{\theta,9}\,\lambda^{(0)}\,\gamma_{\theta,9}^{-1} & \lambda^{(0)} &= -\gamma_{\Omega,9}\,\lambda^{(0)T}\,\gamma_{\Omega,9}^{-1} \\ \lambda^{(a)} &= e^{2\pi i v_a}\,\gamma_{\theta,9}\,\lambda^{(a)}\,\gamma_{\theta,9}^{-1} & \lambda^{(a)} &= -\gamma_{\Omega,9}\,\lambda^{(a)T}\,\gamma_{\Omega,9}^{-1}\,. \end{aligned} \tag{2.7}$$

In the R sector, states are labeled by weights $|\pm\frac{1}{2}, \pm\frac{1}{2}, \pm\frac{1}{2}, \pm\frac{1}{2}\rangle$ of a spinor representation of SO(8), with an odd number of $-1/2$ entries due to the GSO projection. We obtain four left-handed spacetime spinors, whose Chan-Paton factors $\lambda^{(0)}$, $\lambda^{(a)}$ suffer the projections (2.7). Right-handed fermions suffer the opposite projection and provide their antiparticles. Bosonic and fermionic states form multiplets of $\mathcal{N} = 1$ supersymmetry.

[2]Following the classification introduced in [32] for six-dimensional models.

**The 55 sector** The structure of the 55 sector ressembles that of the 99. The only difference arises in the $\Omega$ projection, which acts with $+1$ sign on states associated to complex planes with Dirichlet boundary conditions. The projection for $D5_3$-branes are given by

$$\begin{aligned} \lambda^{(0)} &= \gamma_{\theta,5}\, \lambda^{(0)}\, \gamma_{\theta,5}^{-1} & \lambda^{(0)} &= -\gamma_{\Omega,5}\, \lambda^{(0)T}\, \gamma_{\Omega,5}^{-1} \\ \lambda^{(a)} &= e^{2\pi i v_a}\, \gamma_{\theta,5}\, \lambda^{(a)}\, \gamma_{\theta,5}^{-1} & \lambda^{(a)} &= \pm\gamma_{\Omega,5}\, \lambda^{(a)T}\, \gamma_{\Omega,5}^{-1}\,, \end{aligned} \tag{2.8}$$

with positive sign for $a = 1, 2$ and negative for $a = 3$. The projections being identical for bosonic and fermionic states, they form $\mathcal{N} = 1$ supermultiplets.

**The 59+95 sector** In the NS sector there are zero modes along the directions with DN boundary conditions, hence, states are labeled by an internal SO(4) spinor weight $|s_1, s_2\rangle$, with $s_i = \pm 1/2$. The GSO projection requires $s_1 = s_2$. The orbifold projection imposes the following projection on the Chan-Paton factors of 59 and 95 states

$$\lambda_{59} = e^{\pm 2\pi i (v_1+v_2)/2}\, \gamma_{\theta,5}\, \lambda_{59}\, \gamma_{\theta,9}^{-1} \qquad \lambda_{95} = e^{\pm 2\pi i (v_1+v_2)/2}\, \gamma_{\theta,9}\, \lambda_{95}\, \gamma_{\theta,5}^{-1}\,, \tag{2.9}$$

with the positive and negative signs correspond to the states $|\frac{1}{2}, \frac{1}{2}\,\rangle$ and $|-\frac{1}{2}, -\frac{1}{2}\,\rangle$, respectively. The 59 and 95 sectors are related by the action of $\Omega$ so in fact determining the spectrum in one of them is enough.

The zero modes in the R sector arise along the directions with NN boundary conditions. States are labeled by an SO(4) spinor weight $|s_3; s_0\rangle$, where $s_0$ determines the spacetime fermion chirality. The GSO projection requires $s_0 = s_3$. The projection for these states is

$$\lambda_{59} = e^{\pm 2\pi i v_3/2}\, \gamma_{\theta,5}\, \lambda_{59}\, \gamma_{\theta,9}^{-1} \qquad \lambda_{95} = e^{\pm 2\pi i v_3/2}\, \gamma_{\theta,9}\, \lambda_{95}\, \gamma_{\theta,5}^{-1}\,, \tag{2.10}$$

with the positive sign for $|\frac{1}{2}; \frac{1}{2}\,\rangle$ and the negative for $|-\frac{1}{2}; -\frac{1}{2}\rangle$. As before, $\Omega$ relates the states in the 59 and 95 sectors. Notice also that since $\sum_{a=1}^{3} v_a = 0$, the states form $\mathcal{N} = 1$ supersymmetric multiplets.

**The $\bar{9}\bar{9}$, $\bar{5}\bar{5}$, and $\bar{5}\bar{9} + \bar{9}\bar{5}$ sectors.** Let us move on to the projection for open string states in antibrane-antibrane sectors. In these sectors the GSO projection is the same as in brane-brane sectors. The only difference with respect to the 55, 99 and 59+95 sectors arises in the action of the orientifold projection on the R states, which has an additional $(-1)$ factor. The orientifold projection on the NS states remains unchanged. This fact has been discussed in [24], and reflects the fact that the orientifold projection distinguishes branes from antibranes. In particular, it respects the supersymmetries unbroken by the former and broken by the latter. Notice that this sign flip has no net effect on the spectrum in the $\bar{5}\bar{9}$, $\bar{9}\bar{5}$ sectors, since they are not fixed under $\Omega$.

Finally, let us discuss the mixed brane-antibrane and antibrane-brane sectors. As compared with the analogous brane-brane (or antibrane-antibrane) sectors, they have the opposite GSO projection.

**The $9\bar{9}+\bar{9}9$ sectors.** In the NS sector the tachyon $|0,ab\rangle(\lambda_t)_{ab}$ survives the GSO projection, whereas the would-be massless states $\Psi_{-\frac{1}{2}}|0\rangle$ do not. The orbifold projection on the Chan-Paton factors for the tachyons is

$$\lambda_{t,9\bar{9}} = \gamma_{\theta,9}\,\lambda_{t,9\bar{9}}\,\gamma_{\theta,\bar{9}}^{-1} \qquad \lambda_{t,\bar{9}9} = \gamma_{\theta,\bar{9}}\,\lambda_{t,\bar{9}9}\,\gamma_{\theta,9}^{-1}\,. \tag{2.11}$$

Since $\Omega$ relates the $9\bar{9}$ and $\bar{9}9$ sectors, the tachyons are real fields. The conditions under which the orientifolds are free of tachyons will be discussed below.

In the R sector, the GSO projection selects SO(8) weight vectors $|\pm\frac{1}{2},\pm\frac{1}{2},\pm\frac{1}{2},\pm\frac{1}{2}\rangle$ with an even number of $-1/2$ entries. From these states we get right-handed spacetime fermions with Chan-Paton factors $\lambda^{(0)}$, $\lambda^{(a)}$ constrained by the projections

$$\begin{aligned} \lambda_{9\bar{9}}^{(0)} &= \gamma_{\theta,9}\,\lambda_{9\bar{9}}^{(0)}\,\gamma_{\theta,\bar{9}}^{-1} & \qquad \lambda_{\bar{9}9}^{(0)} &= \gamma_{\theta,\bar{9}}\,\lambda_{\bar{9}9}^{(0)}\,\gamma_{\theta,9}^{-1} \\ \lambda_{9\bar{9}}^{(a)} &= e^{2\pi i v_a}\,\gamma_{\theta,9}\,\lambda_{9\bar{9}}^{(a)}\,\gamma_{\theta,\bar{9}}^{-1} & \qquad \lambda_{\bar{9}9}^{(a)} &= e^{2\pi i v_a}\,\gamma_{\theta,\bar{9}}\,\lambda_{\bar{9}9}^{(a)}\,\gamma_{\theta,9}^{-1}\,. \end{aligned} \tag{2.12}$$

Actually, $\Omega$ relates states in the $9\bar{9}$ and $\bar{9}9$, so it is enough to compute only one of these sectors.

**The $5\bar{5}+\bar{5}5$ sector.** The spectrum in this sector is completely analogous to that in the $9\bar{9}+\bar{9}9$ sector. The only difference, as we may recall from the comparison of 55 and 99 sectors, might arise in the projection imposed by $\Omega$ in directions with D boundary conditions. However, since $\Omega$ does not impose any projection on either the $5\bar{5}$ or the $\bar{5}5$ sectors (rather it maps one to the other), no such difference arises at the level of the spectrum.

**The $9\bar{5}+\bar{5}9$ and $\bar{9}5+5\bar{9}$ sectors.** These sectors are analogous to the $59+95$ sector, with a few modifications due only to the opposite GSO projections. The NS states are labeled by a weight vector $|s_1,s_2\rangle$ with respect to the SO(4) corresponding to the DN directions. The GSO projection in this case selects $s_1=-s_2$. The projections on the Chan-Paton factors of these states in the $9\bar{5}+\bar{5}9$ sector are

$$\lambda_{9\bar{5}} = e^{\pm 2\pi i(v_1-v_2)/2}\,\gamma_{\theta,9}\,\lambda_{9\bar{5}}\,\gamma_{\theta,\bar{5}}^{-1} \qquad \lambda_{\bar{5}9} = e^{\pm 2\pi i(v_1-v_2)/2}\,\gamma_{\theta,\bar{5}}\,\lambda_{\bar{5}9}\,\gamma_{\theta,9}^{-1}\,, \tag{2.13}$$

with the positive and negative signs for the states $|\frac{1}{2},-\frac{1}{2}\,\rangle$ and $|\frac{1}{2},-\frac{1}{2}\rangle$ respectively. The $\Omega$ projection relates the $9\bar{5}$ and the $\bar{5}9$ sectors. The projections in the $\bar{9}5+5\bar{9}$ sector are completely analogous.

States in the R sector are labelled by an SO(4) spinor weight $|s_3;s_0\rangle$, with $s_3=-s_0$ due to the GSO projection, and where $s_0$ defines the spacetime chirality. The projection on the Chan-Paton factors in the $9\bar{5}+\bar{5}9$ sector is

$$\begin{aligned} \lambda_{9\bar{5}} &= e^{\pm 2\pi i v_3/2}\gamma_{\theta,9}\lambda_{9\bar{5}}\gamma_{\theta,\bar{5}}^{-1} \\ \lambda_{\bar{5}9} &= e^{\pm 2\pi i v_3/2}\gamma_{\theta,\bar{5}}\lambda_{\bar{5}9}\gamma_{\theta,9}^{-1}\,, \end{aligned} \tag{2.14}$$

| Sector | Gauge bosons | Scalar fields | Fermion$_+$ | Fermion$_-$ |
|---|---|---|---|---|
| $99^*$ | $\prod_{i=1}^N U(v_i)$ | $\sum_{i=1}^N (v_i, \overline{v}_{i+\ell_a})$ | $\sum_{i=1}^N (v_i, \overline{v}_{i+\ell_a})$ | $\sum_{i=1}^N \mathrm{Adj}_i$ |
| $\overline{99}^*$ | $\prod_{i=1}^N U(w_i)$ | $\sum_{i=1}^N (w_i, \overline{w}_{i+\ell_a})$ | $\sum_{i=1}^N (w_i, \overline{w}_{i+\ell_a})$ | $\sum_{i=1}^N \mathrm{Adj}_i$ |
| $9\overline{9}$ | | $\sum_{i=1}^N (v_i, \overline{w}_i)$ (tachyons) | $\sum_{i=1}^N (v_i, \overline{w}_i)$ | $\sum_{i=1}^N (v_i, \overline{w}_{i+\ell_a})$ |
| $55^*$ | $\prod_{i=1}^N U(n_i)$ | $\sum_{i=1}^N (n_i, \overline{n}_{i+\ell_a})$ | $\sum_{i=1}^N (n_i, \overline{n}_{i+\ell_a})$ | $\sum_{i=1}^N \mathrm{Adj}_i$ |
| $\overline{55}^*$ | $\prod_{i=1}^N U(m_i)$ | $\sum_{i=1}^N (m_i, \overline{m}_{i+\ell_a})$ | $\sum_{i=1}^N (m_i, \overline{m}_{i+\ell_a})$ | $\sum_{i=1}^N \mathrm{Adj}_i$ |
| $5\overline{5}$ | | $\sum_{i=1}^N (n_i, \overline{m}_i)$ (tachyons) | $\sum_{i=1}^N (n_i, \overline{m}_i)$ | $\sum_{i=1}^N (n_i, \overline{m}_{i+\ell_a})$ |
| $59$ | | $\sum_{i=1}^N \left(n_i, \overline{v}_{i-\frac{1}{2}\ell_3}\right)$ | $\sum_{i=1}^N \left(n_i, \overline{v}_{i-\frac{1}{2}\ell_3}\right)$ | |
| $\overline{5}\overline{9}$ | | $\sum_{i=1}^N \left(m_i, \overline{w}_{i-\frac{1}{2}\ell_3}\right)$ | $\sum_{i=1}^N \left(m_i, \overline{w}_{i-\frac{1}{2}\ell_3}\right)$ | |
| $5\overline{9}$ | | $\sum_{i=1}^N \left(n_i, \overline{w}_{i+\frac{1}{2}(\ell_1-\ell_2)}\right)$ | | $\sum_{i=1}^N \left(n_i, \overline{w}_{i-\frac{1}{2}\ell_3}\right)$ |
| $\overline{5}9$ | | $\sum_{i=1}^N \left(m_i, \overline{v}_{i+\frac{1}{2}(\ell_1-\ell_2)}\right)$ | | $\sum_{i=1}^N \left(m_i, \overline{v}_{i+\frac{1}{2}\ell_3}\right)$ |

**Table 1:** The table shows the spectrum obtained in the different open string sectors after imposing the orbifold projections, but before the $\Omega$ projection. To obtain the final spectrum, the orientifold projection must be imposed on the sectors indicated by an asterisk. In this last step it must be kept in mind that in antibrane-antibrane sectors the R states get an additional $(-1)$ factor under the action of $\Omega$. Scalar fields are complex, save for the tachyon fields which are real.

with the positive and negative signs for the states $|\frac{1}{2}; -\frac{1}{2}\rangle$ and $|-\frac{1}{2}; \frac{1}{2}\rangle$, respectively. As usual $\Omega$ relates the $9\bar{5}$ and $\bar{5}9$. The projections for the $\bar{9}5 + 5\bar{9}$ sector are obtained analogously.

In order to illustrate the kind of states survive the projections, let us consider Chan-Paton matrices of the form

$$\gamma_{\theta,9} = \mathrm{diag}\left(1_{v_0}, e^{2\pi i \frac{1}{N}} 1_{v_1}, \ldots, e^{2\pi i \frac{N-1}{N}} 1_{v_{N-1}}\right) \tag{2.15}$$

and analogous expressions for $\overline{\mathrm{D9}}$-, D5- and $\overline{\mathrm{D5}}$-branes, with the numbers of entries replaced by $w_i$, $n_i$ and $m_i$, respectively. We must clarify, however, that our projection rules do not depend on the specific form of these matrices. Indeed, in some of our models the Chan-Paton matrices will differ from the expression above.

The states resulting from the *orbifold* projection with such matrices are shown in table 1. The subindices denote the gauge group under which the state transforms. They are defined mod N, hence for negative indices we have $v_{-i} = v_{N-i}$. The choice of fermions chiralities is mainly the usual one for supersymmetric sectors. In brane-antibrane sectors, the opposite chirality arises from the opposite GSO projection. The orientifold projection will impose additional conditions on these states in some of the sectors, indicated with an asterisk. For sectors exchanged by the orientifold projection only one of them is shown. For these states the orientifold projection does not impose additional constraints.

One important point to be stressed is that tachyons only appear when the model contains coincident (or very close) branes of the same type and with identical Chan-Paton phases. Those tachyons obviously signal instabilities due to brane-antibrane annihilation. Therefore the construction of type IIB orientifolds with branes and antibranes but without tachyons is fairly simple. It only requires that branes and antibranes sit at different points in the internal space (or have different Wilson lines, in a T-dual picture), or if they are coincident, that their Chan-Paton matrices project out the tachyons (that is $n_i = 0$ or $m_i = 0$, and $v_i = 0$ or $w_i = 0$, for all $i = 1, \ldots, N$). In sections 3 and 4 we give explicit examples of tachyon-free models, and comment on some of their properties.

### 2.3 Cylinder partition function

In this section we indicate how the open string spectrum derived above shows up in the cylinder partition functions (see also [24, 27]). We include it here in order to provide a cross check for the correctness of the construction in section 2.1, and also in order to illustrate how different signs must be suitable flipped, in a concrete amplitude, in order to take the presence of anti-branes into account. We also show how RR tadpole cancellation is achieved by performing the usual transformation to the closed string channel. We restrict our analysis to configurations containing D9-branes and $\overline{\text{D9}}$-branes. Other amplitudes can be similarly analyzed.

The number of states at each mass level $M$ can be obtained by looking at the multiplicity of the term $q^{M^2}$ ($q = e^{-2\pi t}$) in the partition function. We will be interested in negative powers, indicating the presence of a potential tachyon, and in massless fields.

Consider first an array of D$p$ and D$q$-branes. The cylinder amplitudes, in $D$ dimensions are given by (see for instance [9])

$$\mathcal{C}_{pq} = \frac{V_D}{2N} \sum_{k=0}^{N-1} \int_0^\infty \frac{dt}{t} \, (8\pi^2 \alpha' t)^{-D/2} \mathcal{Z}_{pq}(\theta^k) \, , \tag{2.16}$$

where

$$\mathcal{Z}_{pq}(\theta^k) = \frac{1}{2} \operatorname*{Tr}_{pq} \left\{ (1 + (-1)^F) \theta^k \, \mathrm{e}^{-2\pi t L_0} \right\} . \tag{2.17}$$

The trace is over open string states ending at the corresponding brane. For instance in the **99** sector, where boundary conditions are NN in all directions we have

$$\mathcal{Z}_{99}(\theta^k) = \frac{1}{2} \sum_{\alpha,\beta=0,1/2} \eta_{\alpha,\beta} Z\Big[{\alpha \atop \beta}\Big] \, (\mathrm{Tr}\, \gamma_{k,9} \, \mathrm{Tr}\, \gamma_{k,9}^{-1}) \, , \tag{2.18}$$

with $\eta_{1/2,0} = \eta_{0,1/2} = -\eta_{0,0} = -1$.

Here we have defined

$$Z[{\alpha \atop \beta}] = \left[\frac{\vartheta[{\alpha \atop \beta}]}{\eta^3}\right]^{\frac{1}{2}(D-2)} \prod_{a=1}^{\frac{1}{2}(10-D)} \frac{\vartheta[{\alpha \atop \beta+k\frac{\ell_a}{N}}]}{\eta} \frac{-2\sin\pi k v_a\,\eta}{\vartheta[{\frac{1}{2} \atop \frac{1}{2}+k\frac{\ell_a}{N}}]}\,. \tag{2.19}$$

We also define, for later convenience,

$$\begin{aligned}
\mathcal{Z}_{99}^{NS} &= \frac{1}{2N}\sum_{k=0}^{N-1}\left[Z[{0 \atop 0}](t) - Z[{0 \atop \frac{1}{2}}](t)\right]\left(\mathrm{Tr}\,\gamma_{k,9}\,\mathrm{Tr}\,\gamma_{k,9}^{-1}\right)\\
\mathcal{Z}_{99}^{R} &= -\frac{1}{2N}\sum_{k=0}^{N-1} Z[{\frac{1}{2} \atop 0}](t)\left(\mathrm{Tr}\,\gamma_{k,9}\,\mathrm{Tr}\,\gamma_{k,9}^{-1}\right),
\end{aligned} \tag{2.20}$$

which are the separate contributions of open string NS bosons and R fermions respectively. Recall also that $Z[{0 \atop \frac{1}{2}}](t)$ correspond to RR fields in the closed string sector.

In order to extract the leading behaviour for $q = e^{-2\pi t}$ we use the product form of $\vartheta$ function

$$\frac{\vartheta[{\delta \atop \varphi}]}{\eta} = \mathrm{e}^{2i\pi\delta\varphi}\, q^{\frac{1}{2}\delta^2-\frac{1}{24}} \prod_{n=1}^{\infty}\left(1+q^{n+\delta-\frac{1}{2}}\mathrm{e}^{2i\pi\varphi}\right)\left(1+q^{n-\delta-\frac{1}{2}}\mathrm{e}^{-2i\pi\varphi}\right) \tag{2.21}$$

where the Dedekind $\eta$ function is

$$\eta = q^{\frac{1}{24}} \prod_{n=1}^{\infty}(1-q^n)\,. \tag{2.22}$$

Thus, collecting all terms we finally have, for $\eta_{\alpha,\beta} Z[{\alpha \atop \beta}]$

$$\begin{aligned}
Z[{0 \atop 0}](t) &= q^{-\frac{1}{2}} + \left[D-2+2\sum_{a=1}^{\frac{1}{2}(10-D)}\cos\left(\frac{2\pi k\ell_a}{N}\right)\right] + O(q) + \cdots\\
-Z[{0 \atop \frac{1}{2}}](t) &= -q^{-\frac{1}{2}} + \left[D-2+2\sum_{a=1}^{\frac{1}{2}(10-D)}\cos\left(\frac{2\pi k\ell_a}{N}\right)\right] + O(q) + \cdots\\
-Z[{\frac{1}{2} \atop 0}](t) &= -16\prod_{a=1}^{\frac{1}{2}(10-D)}\cos\left(\frac{\pi k\ell_a}{N}\right)\\
-Z[{\frac{1}{2} \atop \frac{1}{2}}](t) &= 0\,.
\end{aligned} \tag{2.23}$$

We see that the potential tachyon field associated to $q^{-1/2}$ disappears from the spectrum. Also, since $D-2+2\sum_{a=1}^{\frac{1}{2}(10-D)}\cos(2\pi k\ell_a/N) = 8\prod_{a=1}^{\frac{1}{2}(10-D)}\cos(\pi k\ell_a/N)$ (for

$\sum_{a=1}^{\frac{1}{2}(10-D)} \ell_a = 0$) bosons and fermions exactly cancel in $\mathcal{Z}_{99}$ at this order (and to all orders due to Riemann identities [28]), exhibiting supersymmetry.

The traces of the twist matrices defined in (2.15) are

$$\mathrm{Tr}\,\gamma_{k,9} = \sum_{j=0}^{N-1} e^{2i\pi kj/N} v_j \tag{2.24}$$

and similarly for $\mathrm{Tr}\,\gamma_{k,9}^{-1}$. Hence, by introducing these expressions in (2.20) and by performing the sum over $k$ we finally obtain

$$\mathcal{Z}_{99}^{NS} = \sum_{i=0}^{N-1} \sum_{a=1}^{\frac{1}{2}(10-D)} \left[(D-2)v_i v_i + v_i v_{i+\ell_a} + v_{i+\ell_a} v_i\right] + O(q) + \cdots . \tag{2.25}$$

with the convention $v_i = v_{N+i}$. These are the exact multiplicities for the gauge bosons $\prod_{i=1}^{N} U(v_i)$ (the factor $D-2$ corresponding to the number of transverse polarizations), and the scalars computed in section 2.2. Recall that an extra 1/2 factor must be included when the orientifold projection is performed. Thus a multiplicity $\frac{1}{2}v_i^2$ would appear. This completes to $\frac{1}{2}v_i(v_i-1)$ multiplicity of $\mathrm{SO}(v_i)$ group when the unoriented string amplitudes are included.

For fermions we have $\mathcal{Z}_{99}^{R} = -\mathcal{Z}_{99}^{NS}$. Multiplicities are now interpreted as coming from positive (internal) chirality fermions and negative chirality adjoint fermions as shown in table 1.

Multiplicities in $\overline{9}\overline{9}$ sector are obtained by simply replacing $v_i \to w_i$ above.

In order to obtain the states multiplicities in $\overline{9}9 + 9\overline{9}$ cylinders we must consistently change the sign of the closed sector RR charges. This amounts to flipping the sign of $Z[{0 \atop \frac{1}{2}}](t)$ above. Namely

$$\begin{aligned}
\mathcal{Z}_{9\overline{9}+\overline{9}9}^{NS} &= \frac{1}{2N} \sum_{k=0}^{N-1} \left[Z[{0 \atop 0}](t) + Z[{0 \atop \frac{1}{2}}](t)\right] \left(\mathrm{Tr}\,\gamma_{k,\overline{9}}\,\mathrm{Tr}\,\gamma_{k,9}^{-1} + \mathrm{Tr}\,\gamma_{k,9}\,\mathrm{Tr}\,\gamma_{k,\overline{9}}^{-1}\right) \\
\mathcal{Z}_{9\overline{9}+\overline{9}9}^{R} &= -\frac{1}{2N} \sum_{k=0}^{N-1} \left[Z[{\frac{1}{2} \atop 0}](t)\right] \left(\mathrm{Tr}\,\gamma_{k,\overline{9}}\,\mathrm{Tr}\,\gamma_{k,9}^{-1} + \mathrm{Tr}\,\gamma_{k,9}\,\mathrm{Tr}\,\gamma_{k,\overline{9}}^{-1}\right) .
\end{aligned} \tag{2.26}$$

Hence, we notice, by recalling the expansions (2.23), that NS tadpoles do not cancel. In fact, we find

$$\mathcal{Z}_{9\overline{9}+\overline{9}9}^{NS} = q^{-\frac{1}{2}} \sum_{j=0}^{N-1} 2 v_j w_j + O(q) + \cdots , \tag{2.27}$$

which corresponds to $\overline{9}9$ sector tachyons $\sum_{i=0}^{N-1}[(v_i, \overline{w}_i) + (\overline{v}_i, w_i)]$.

Finally for fermions we obtain

$$\mathcal{Z}^R_{9\bar{9}+\bar{9}9} = 2 2^{\frac{D-4}{2}} \left[ \sum_{i=0}^{N-1} 2 v_i w_i + \sum_{a=1}^{\frac{1}{2}(10-D)} v_i w_{i+\ell_a} + v_{i+\ell_a} w_i \right] + O(q) + \cdots . \tag{2.28}$$

Again, these are the multiplicities which correspond to massless fermionic states obtained above.

**RR tadpole cancellation** Cancellation of RR tadpoles is a basic requirement for the consistency of the theory. In type IIB theory with only D-branes (no anti-branes), tadpole cancellation is equivalent to anomaly cancellation in $D = 10, 6$ dimensions, however, it is generally stronger in $D = 4$ [29]. Also, if RR tadpoles are absent so must be NS tadpoles due to supersymmetry. This last result is thus not expected in models with anti-branes where supersymmetry is broken. NS tadpoles should manifest as a background redefinition [30].

In order to analyze the tadpole divergences it is useful to rewrite the partition function as

$$\begin{aligned} \mathcal{Z}^{NS} = -\frac{1}{2N} \sum_{k=0}^{N-1} \Big\{ & Z[{}^{0}_{\frac{1}{2}}](t) \left( \mathrm{Tr}\, \gamma_{k,9} - \mathrm{Tr}\, \gamma_{k,\bar{9}} \right) \left( \mathrm{Tr}\, \gamma^{-1}_{k,9} - \mathrm{Tr}\, \gamma^{-1}_{k,\bar{9}} \right) + \\ & + Z[{}^{0}_{0}](t) \left( \mathrm{Tr}\, \gamma_{k,9} + \mathrm{Tr}\, \gamma_{k,\bar{9}} \right) \left( \mathrm{Tr}\, \gamma^{-1}_{k,9} + \mathrm{Tr}\, \gamma^{-1}_{k,\bar{9}} \right) \Big\} \end{aligned} \tag{2.29}$$

and

$$\mathcal{Z}^{R} = -\frac{1}{2N} \sum_{k=0}^{N-1} Z[{}^{0}_{0}](t) \left( \mathrm{Tr}\, \gamma_{k,9} + \mathrm{Tr}\, \gamma_{k,\bar{9}} \right) \left( \mathrm{Tr}\, \gamma^{-1}_{k,9} + \mathrm{Tr}\, \gamma^{-1}_{k,\bar{9}} \right) , \tag{2.30}$$

where we have used (2.20) (and a similar expression with $9 \to \bar{9}$ for the $\overline{99}$ sector) and (2.26).

RR tadpoles divergences are contained in $Z[{}^{0}_{\frac{1}{2}}](t)$ and have been analyzed in different situations in type IIB. We rederive them here for completeness. In fact, by performing a modular transformation $t \to 1/t$ and looking at small values of $t$ it can easily obtain

$$Z[{}^{0}_{\frac{1}{2}}](1/t) = t^{\frac{1}{2}(D-2)} \prod_{a=1}^{\frac{1}{2}(10-D)} \left( -2 \sin \frac{\pi k \ell_a}{N} \right) . \tag{2.31}$$

Nevertheless we note that, regarding tadpole cancellation, the net effect of the introduction of anti D9-branes, with respect to type IIB theory with only D9-branes, is that we must replace

$$\mathrm{Tr}\, \gamma_{k,9} \to \mathrm{Tr}\, \gamma_{k,9} - \mathrm{Tr}\, \gamma_{k,\bar{9}} \tag{2.32}$$

since we must require the coefficient of $Z[{}^{0}_{\frac{1}{2}}]$ in (2.29) to vanish. The rule (2.32) reflects the fact that branes and antibranes have opposite charges with respect to the RR fields.

For instance, in a six-dimensional type IIB orbifold compactification (non orientifold yet), the only tadpoles arise from cylinder amplitudes. The corresponding equations can be extracted from [8] and lead to

$$\mathrm{Tr}\,\gamma_{k,9} - \mathrm{Tr}\,\gamma_{k,\overline{9}} - 4\sin^2\frac{\pi k}{N}(\mathrm{Tr}\,\gamma_{k,5,L} - \mathrm{Tr}\,\gamma_{k,\overline{5},L}) = 0 \tag{2.33}$$

for $5, \overline{5}$-branes at orbifold fixed point $L$.

Notice that by computing the traces as in (2.24), these equations are equivalent to (we have dropped the $L$ index for $m_i, n_i$ here)

$$-2n_i + n_{i+1} + n_{i-1} + v_i + 2m_i - m_{i+1} - m_{i-1} = 0\,. \tag{2.34}$$

Precisely, these are the conditions for cancellation of anomalies of the six dimensional gauge theories living on branes and anti-branes with the spectrum computed in section 2.1 (before the orientifold projection).

It is also possible to show that the Green-Schwarz mechanism that cancels the residual anomalies in six [31, 32] and four [33] dimensions works in models with antibranes. In this respect, it is important to notice that the fermions in brane-antibrane sectors have chirality opposite to that of usual matter fermions in supersymmetric models. Hence their contribution to the anomaly has the opposite sign as well. This sign also emerges in the GS counterterms because the coupling of antibranes to RR fields is opposite to that of branes. Hence cancellation follows in the usual way. For antibrane-antibrane sectors, fermions have the usual chirality, and the GS countertems also have the usual sign (since the $-1$ signs of the two antibrane couplings give an overall $+1$).

## 3. Some six-dimensional examples

In this section we construct several explicit non-supersymmetric type IIB orientifolds in six dimensions with chiral spectrum. These models are free of tachyons, and illustrate a number of interesting generic features of compactifications with branes and antibranes. Also, since chiral theories are potentially anomalous in six dimensions, the cancellation of anomalies in the models we present serves as a useful check of the rules proposed to construct the spectrum in section 2.

### 3.1 A non-supersymmetric $Z_3$ model

A simple possibility to obtain non-supersymmetric models is to modify slightly one of the familiar supersymmetric compactifications by the introduction of antibranes, in a way consistent with cancellation of RR tadpoles. In order to illustrate the basic idea, let us consider a simple model related to the six-dimensional $\mathbb{Z}_3$ model in [8]. The construction will be mostly intuitive in terms of a T-dual version, where $T^4/\mathbb{Z}_3$ is modded out by $\Omega R_1 R_2$, and the model contains no D9-branes. It contains D5-branes (and $\overline{\text{D5}}$-branes) sitting at points in the compact space. As discussed in section 2.2, the twisted tadpole cancellation conditions for the branes at the origin are obtained from those in [8] by replacing $\operatorname{Tr}\gamma_{\theta^k,5} \to \operatorname{Tr}\gamma_{\theta^k,5} - \operatorname{Tr}\gamma_{\theta^k\bar{5}}$, so we have

$$
\begin{aligned}
\operatorname{Tr}\gamma_{\theta,5} - \operatorname{Tr}\gamma_{\theta,\bar{5}} &= 8 \\
Tr\gamma_{\theta^2,5} - \operatorname{Tr}\gamma_{\theta^2,\bar{5}} &= 8\,.
\end{aligned}
\tag{3.1}
$$

These can be satisfied, for instance, by choosing

$$
\gamma_{\theta,5} = \mathbb{1}_4\,; \qquad \gamma_{\theta,\bar{5}} = \operatorname{diag}\left(e^{2\pi i\frac{1}{3}}1_4, e^{2\pi i\frac{2}{3}}1_4\right). \tag{3.2}
$$

The matrices $\gamma_{\Omega R_1 R_2}$ simply exchanges opposite phases in $\gamma_\theta$. The untwisted tadpole requires a net D5-brane number of 32, $n_5 - n_{\bar{5}} = 32$, and so the model must contain additional D5-branes. Let us consider placing 18 D5-branes at one of the orbifold (rather than orientifold) points fixed under $\mathbb{Z}_3$, and another set of 18 at its image under $\Omega R_1 R_2$. For an orbifold point, twisted tadpole conditions amount to tracelessness of the relevant Chan-Paton matrices. Therefore these D5-branes are described by

$$
\gamma_{\theta,5} = \operatorname{diag}\left(1_6, e^{2\pi i\frac{1}{3}}1_6, e^{2\pi i\frac{2}{3}}1_6\right). \tag{3.3}
$$

The gauge group arising from branes (and antibranes) at the origin is $\mathrm{SO}(4)_{55} \times \mathrm{U}(4)_{\bar{5}\bar{5}}$, with non-supersymmetric matter content given in the following table.[3]

| Sector | Complex scalars | Fermion$_+$ | Fermion$_-$ |
|---|---|---|---|
| $55$ | | | $(\mathrm{Adj}, 1)$ |
| $\bar{5}\bar{5}$ | $(1, \begin{smallmatrix}\square\\ \square\end{smallmatrix}) + \left(1, \overline{\begin{smallmatrix}\square\\ \square\end{smallmatrix}}\right)$ | $(1, \square\!\square)$ | $(1, \mathrm{Adj})$ |
| $5\bar{5} + \bar{5}5$ | | | $(\square, \square)$ |

The D5-branes sitting at the orbifold fixed points give a $D = 6$, $\mathcal{N} = 1$ spectrum with group $\mathrm{U}(6)^3$, and hypermultiplets in $(\square, \overline{\square}, 1) + (1, \square, \overline{\square}) + (\overline{\square}, 1, \square)$.

Since the orientifold projection on the closed string sector is exactly that in [8], it gives rise to the $D = 6$, $\mathcal{N} = 1$ supergravity and dilaton multiplets, and the following set of $D = 6$, $\mathcal{N} = 1$ matter multiplets: two hypermultiplets in the untwisted sector, and nine hyper- and nine tensor multiplets in the twisted sectors. It is easy to check that both gauge and gravitational irreducible anomalies cancel in the model. As usual, the residual non-abelian and U(1) anomalies are cancelled by the GS mechanism mediated by the tensor and hypermultiplets [31, 32].

[3] Notice that the projection with $\Omega R_1 R_2$ has some additional $(-1)$ signs as compared with the $\Omega$ projection.

This theory can be thought of as a toy model where we have a supersymmetric sector of branes and a non-supersymmetric sector containing antibranes. Supersymmetry breaking would be transmitted from the latter to the former by exchange of bulk fields. However, the model suffers from the following serious drawback. As discussed in a number of papers (see e.g. [34]), the lack of BPS properties in brane-antibrane systems leads to an attractive force between the two kinds of objects. In brane-brane (or antibrane-antibrane) systems, the repulsive force mediated by exchange of RR closed string fields is exactly cancelled by that from exchange of NS-NS fields, leading to no net force. In brane-antibranes systems, however, the RR piece becomes attractive and branes and antibranes tend to come close to each other. Once the distance reaches a critical value, a tachyon develops, signaling the possibility of annihilation and decay to the vacuum. This type of phenomenon takes place in the model we have constructed. Indeed, the D5-branes sitting at the $\mathbb{Z}_3$ orbifold point are not stuck to it, and are free to move into the bulk as a dynamical D5-brane. This transition very likely takes place due to the attractive force induced by the D5-branes located (and actually, stuck) at the origin. Once in the bulk, the D5-branes approach the origin, tachyons develop, branes and antibranes annihilate and the theory decays to the supersymmetric model in [8], where the Chan-Paton matrices for D5-branes are

$$\gamma_{\theta,5} = \mathrm{diag}\left(1_{16}, e^{2\pi i\frac{1}{3}}1_8, e^{2\pi i\frac{2}{3}}1_8\right). \tag{3.4}$$

Therefore, the initial mode should be thought of as an excited non-supersymmetric state of this supersymmetric vacuum. It is straightforward to construct this type of non-supersymmetric compactifications, by nucleating brane-antibrane pairs in any supersymmetric model, and separating the members of these pairs. Since this type of models is not really new or interesting for our purposes, the models we construct in the following will be safe against this type of decay. This can be achieved, as we show in section 4 by considering branes and antibranes stuck at different fixed points. A different possibility is to consider branes and antibranes with different world-volume dimension (since tachyons only appear in $p\overline{p}$ sectors). Models of this kind are presented in next subsection.

### 3.2 Non-supersymmetric models with vector structure

In this section we construct a different kind of models, quite unrelated to supersymmetric orientifolds. They illustrate the fact that the introduction of antibranes sometimes allows to satisfy RR tadpole cancellation conditions which could not be satisfied in supersymmetric models containing only D-branes.

The models we consider are obtained by modding out type IIB theory on $T^4/\mathbb{Z}_N$ (with $N = 2, 4, 6$) by an orientifold projection $\Omega$ that preserves $D = 6$, $\mathcal{N} = 1$ supersymmetry in the closed string spectrum. The orientifold projection we are interested in discussing, however, differs from that in [8] and was first discussed in [35]. Concretely, its action on the order two twisted sector of $\mathbb{Z}_N$ is such that the RR states are left-right symmetric (rather than antisymmetric) combinations. In other words, the orientifold projection gives rise to a tensor multiplet (rather than a hypermultiplet) in the $\mathbb{Z}_2$ twisted sectors. This projection imposes a constraint on the Chan-Paton matrices of the corresponding twist [35]

$$\gamma_\Omega = \gamma_{\theta^{N/2}}\gamma_\Omega\gamma_{\theta^{N/2}}^T \tag{3.5}$$

for D9- and D5-branes (and also for antibranes). The constraint implies the bundles on the D-branes have vector structure (see [32] for a discussion of vector structure in orientifold models). This type of projection was studied in [36, 37] in the context of orientifolds of $\mathbb{C}^2/\mathbb{Z}_N$, which provide a local description of the fixed point of the compact models.

It is not possible to construct a supersymmetric orientifold of $T^4/\mathbb{Z}_N$ where all fixed points are of this type.[4] The reason is that the cancellation of the untwisted tadpole generated by the orientifold planes requires $n_5 = -32$. Clearly it is not possible to cancel this tadpole by introducing D5-branes. However, as we show below, it is certainly possible to construct consistent tachyon-free non-supersymmetric models, where this tadpole is neatly cancelled by the introduction of 32 $\overline{\text{D5}}$-branes. Untwisted tadpoles also require $n_9 = 32$, as in the usual case. In these models, tachyons are absent because branes and antibranes have different world-volume dimension, hence no anihilation is possible.

Before entering the detailed construction of the model, let us introduce some conventions, which differ slightly from those in section 2.1. The Chan-Paton matrices for the twist $\gamma_\theta$ will be of the form

$$\gamma_{\theta,9} = \mathrm{diag}\left(1_{v_0}, e^{2\pi i\frac{1}{N}}1_{v_1}, \ldots, e^{2\pi i\frac{N-1}{N}}1_{v_{N-1}}\right) \tag{3.6}$$

(and analogously for $\gamma_{\theta,\bar{5}}$), with $v_i = v_{N-i}$ due to the orientifold symmetry. The matrices $\gamma_{\Omega,9}$ will be of the form

$$\gamma_{\Omega,9} = \begin{pmatrix} 1_{v_0} & & & & \\ & & & & 1_{v_1} \\ & & & \cdots & \\ & & 1_{v_{N/2}} & & \\ & \cdots & & & \\ & 1_{v_{N-1}} & & & \end{pmatrix}, \quad \gamma_{\Omega,\bar{5}} = \begin{pmatrix} \epsilon_{v_0} & & & & \\ & & & & 1_{v_1} \\ & & & \cdots & \\ & & \epsilon_{v_{N/2}} & & \\ & \cdots & & & \\ & -1_{v_{N-1}} & & & \end{pmatrix}.$$

These conventions are useful to make contact with [36, 37], where the twisted tad-

[4] It is however possible to construct mixed models with fixed points of different kinds, see [35] for a $\mathbb{Z}_2$ example.

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poles for these orientifolds were computed. The general expression can be extracted from these references[5]

$$\mathrm{Tr}\,\gamma_{\theta^k,9} + 4\sin^2\frac{\pi k}{N}\,\mathrm{Tr}\,\gamma_{\theta^k,\bar{5}} - 32\delta_{k,0 \bmod 2} = 0 \tag{3.7}$$

for branes at $\mathbb{Z}_N$ fixed points. Since these models are in a sense the simplest six-dimensional compactifications with vector structure, we now turn to the explicit construction of the $\mathbb{Z}_2$, $\mathbb{Z}_4$ and $\mathbb{Z}_6$ models.

#### 3.2.1 The $\mathbb{Z}_2$ model

Consider modding out type IIB on $T^4/\mathbb{Z}_2$ by the orientifold projection $\Omega$ that selects the tensor multiplet in all $\mathbb{Z}_2$ twisted sectors. The closed string spectrum gives a set of $D = 6$, $\mathcal{N} = 1$ multiplets, concretely 4 hypermultiplets in the untwisted sector and 16 tensor multiplets in the twisted sector.

Concerning the open string spectrum, we consider, as discussed above, $n_9 = n_{\bar{5}} = 32$ to satisfy the untwisted tadpoles. We will also consider the model where all the $\overline{\text{D5}}$-branes sit at one $\mathbb{Z}_2$ fixed point. Other models can be constructed analogously. RR tadpoles for such configuration are cancelled by choosing

$$\begin{aligned}\gamma_{\theta,9} &= \mathrm{diag}(1_{16}, -1_{16})\\ \gamma_{\theta,\bar{5}} &= \mathrm{diag}(1_{16}, -1_{16})\,.\end{aligned} \tag{3.8}$$

The gauge group is $[\mathrm{SO}(16)\times\mathrm{SO}(16)]_{99}\times[\mathrm{USp}(16)\times\mathrm{USp}(16)]_{\bar{5}\bar{5}}$. The matter content arising from the projections in section 2 is

$$\begin{array}{lll}
\mathbf{99}\text{ sector}: & \text{Complex Scalars}: & 2(16,16)\\
 & \text{Fermions}_+\ : & (16,16)\\
 & \text{Fermions}_-\ : & (120,1)+(1,120)\\
\bar{\mathbf{5}}\bar{\mathbf{5}}\text{ sector}: & \text{Complex Scalars}: & 2(16,16)\\
 & \text{Fermions}_+\ : & (16,16)\\
 & \text{Fermions}_-\ : & (120,1)+(1,120)\\
\bar{\mathbf{5}}\mathbf{9}+\mathbf{9}\bar{\mathbf{5}}\text{ sector}: & \text{Complex Scalars}: & (16,1;1,16)+(1,16;16,1)\\
 & \text{Fermions}_-\ : & \frac{1}{2}(16,1;16,1)+\frac{1}{2}(1,16;1,16)\,.
\end{array} \tag{3.9}$$

Here the **120** of USp(16) is actually reducible as **119**+**1**. Also, the factor 1/2 in front of some fermion representations implies they have a (symplectic) Majorana-Weyl constraint, which is possible because they transform in pseudoreal representations of the gauge group (these are the fermions appearing in the more familiar half-hypermultiplets of $D = 6$, $\mathcal{N} = 1$ supersymmetric models).

[5] The tadpole conditions in [36] are slightly different due to a different form of the matrices $\gamma_\theta$. Our expression is taken from section 3.2 in [37], with the only modifications of multiplying the crosscap by a factor of 4 (since our models are six-dimensional rather than four-dimensional) and replacing $\mathrm{Tr}\,\gamma_{\theta^k,5}$ by $-\,\mathrm{Tr}\,\gamma_{\theta^k,\bar{5}}$.

It is easy to check that all irreducible gauge and gravitational anomalies cancel in the model above. As usual, we expect the residual non-abelian and U(1) anomalies to be cancelled through a Green-Schwarz mechanism [31].

### 3.2.2 The $\mathbb{Z}_4$ model

Let us discuss the construction of a orientifold of $T^4/\mathbb{Z}_4$ by this orientifold projection. In the closed string sector we obtain a set of $D=6$, $\mathcal{N}=1$ multiplets. Specifically, the untwisted sector provides two hypermultiplets. Concerning the twisted sector, the model contains four $\mathbb{C}^2/\mathbb{Z}_4$ fixed points, each of which contributes with one hyper- and two tensor multiplets, and six $\mathbb{C}^2/\mathbb{Z}_2$, each contributing with one tensor multiplet. Hence we get 6 hyper- and 14 tensor multiplets.

Regarding the open string sector, untwisted tadpole cancellation conditions require $n_9 = n_{\bar{5}} = 32$. It is possible to show that in the $\mathbb{Z}_4$ model tadpole cancellation conditions do not have solutions if all $\overline{\text{D5}}$-branes are located at a single fixed point. However, if we place $\overline{\text{D5}}$-branes at all four $\mathbb{Z}_4$ fixed points, it is possible to solve the tadpole conditions. In what follows we describe a simple solution, even though there are other possibilities.

Let us consider the D9-branes to be described by the following Chan-Paton matrix

$$\gamma_{\theta,9} = \mathrm{diag}\left(1_8, e^{2\pi i\frac{1}{4}}1_8, e^{2\pi i\frac{2}{4}}1_8, e^{2\pi i\frac{3}{4}}1_8\right). \tag{3.10}$$

We also consider a set of 8 $\overline{\text{D5}}$-branes sitting at each of the four $\mathbb{Z}_4$ fixed points labeled by $L=1,\ldots,4$. Their Chan-Paton matrix is

$$\gamma_{\theta,\bar{5},L} = \mathrm{diag}\left(1_4, e^{2\pi i\frac{2}{4}}1_4\right). \tag{3.11}$$

These choices satisfy the tadpole conditions (3.7) at each $\mathbb{Z}_4$ fixed point.

The gauge group in this model is

$$\mathrm{SO}(8)\times\mathrm{U}(8)\times\mathrm{SO}(8)\times\prod_{L=1}^{4}[\mathrm{USp}(4)\times\mathrm{USp}(4)]_L\,. \tag{3.12}$$

The matter content is given by

$$\begin{array}{lll}
\mathbf{99}\text{ sector}: & \text{Complex Scalars}: & (8,\bar{8},1)+(1,8,8)+(8,8,1)+(1,\bar{8},8)\\
 & \text{Fermions}_+\;: & (8,\bar{8},1)+(1,8,8)\\
 & \text{Fermions}_-\;: & (28,1,1)+(1,64,1)+(1,1,28)\\
\mathbf{\bar{5}_L\bar{5}_L}\text{ sector}: & \text{Fermions}_- & (6_L,1)+(1,6_L)\\
\mathbf{\bar{5}_L 9}+\mathbf{9\bar{5}_L}\text{ sector}: & \text{Complex Scalars}: & (4_L,1;1,\bar{8},1)+(1,4_L;1,8,1)\\
 & \text{Fermions}_-\;: & \frac{1}{2}(4_L,1;8,1,1)+\frac{1}{2}(1,4_L;1,1,8)\,,
\end{array} \tag{3.13}$$

where the $\mathbf{6}$ of USp(4) is actually reducible as $\mathbf{5}+\mathbf{1}$. It is easy to see that the complete open and closed string spectrum cancels all gauge and gravitational anomalies.

### 3.2.3 The $\mathbb{Z}_6$ model

Let us consider the six-dimensional $\mathbb{Z}_6$ model, with the orientifold projections selecting the tensor multiplets in $\mathbb{Z}_2$ twisted sectors. The closed string sector gives rise to a set of $D = 6$, $\mathcal{N} = 1$ multiplets. There are two hypermultiplets arising from the untwisted sector. The model contains one $\mathbb{C}^2/\mathbb{Z}_6$ fixed points, which contributes two hyper- and three tensor multiplets, four $\mathbb{C}^2/\mathbb{Z}_3$ orbifold points, contributing one hyper- and one tensor multiplet each, and five $\mathbb{C}^2/\mathbb{Z}_2$ points, contributing one tensor multiplet each. We have a total of 8 hyper- and 12 tensor multiplets.

Concerning the open string spectrum, we have as usual $n_9 = n_{\bar{5}} = 32$. In this case there are solutions to the twisted tadpole cancellation conditions with all $\overline{\mathrm{D5}}$-branes sitting at the origin. The general solution for the Chan-Paton matrices in such configuration is

$$\begin{aligned}\gamma_{\theta,9} &= \mathrm{diag}\left(1_{2K}, e^{2\pi i\frac{1}{6}}1_K, e^{2\pi i\frac{2}{6}}1_{8-K}, e^{2\pi i\frac{3}{6}}1_{16-2K}, e^{2\pi i\frac{4}{6}}1_{8-K}, e^{2\pi i\frac{5}{6}}1_K\right)\\ \gamma_{\theta,\bar{5}} &= \mathrm{diag}\left(1_{16-2K}, e^{2\pi i\frac{1}{6}}1_{8-K}, e^{2\pi i\frac{2}{6}}1_K, e^{2\pi i\frac{3}{6}}1_{2K}, e^{2\pi i\frac{4}{6}}1_K, e^{2\pi i\frac{5}{6}}1_{8-K}\right).\end{aligned} \tag{3.14}$$

A particularly nice solution is obtained for $K = 4$, on which we center for the sake of concreteness, the spectrum for general $K$ being analogous. The gauge group in such case is

$$\mathrm{SO}(8)\times\mathrm{U}(4)\times\mathrm{U}(4)\times\mathrm{SO}(8)\times\mathrm{USp}(8)\times\mathrm{U}(4)\times\mathrm{U}(4)\times\mathrm{USp}(8)\,. \tag{3.15}$$

The matter spectrum is given by

$$\begin{array}{rll}
\mathbf{99}\ \text{sector}: & \text{Complex Scalars}: & (8,\overline{4},1,1)+(1,4,\overline{4},1)+(1,1,4,8)+\\
& & \quad +(8,4,1,1)+(1,\overline{4},4,1)+(1,1,\overline{4},8)\\
& \text{Fermions}_+\ : & (8,\overline{4},1,1)+(1,4,\overline{4},1)+(1,1,4,8)\\
& \text{Fermions}_-\ : & (28,1,1,1)+(1,16,1,1)+\\
& & \quad +(1,1,16,1)+(1,1,1,28)\\
\overline{\mathbf{55}}\ \text{sector}: & \text{Complex Scalars}: & (8,\overline{4},1,1)+(1,4,\overline{4},1)+(1,1,4,8)+\\
& & \quad +(8,4,1,1)+(1,\overline{4},4,1)+(1,1,\overline{4},8)\\
& \text{Fermions}_+\ : & (8,\overline{4},1,1)+(1,4,\overline{4},1)+(1,1,4,8)\\
& \text{Fermions}_-\ : & (28,1,1,1)+(1,16,1,1)+\\
& & \quad +(1,1,16,1)+(1,1,1,28)\\
\overline{\mathbf{5}}\mathbf{9}+\mathbf{9}\overline{\mathbf{5}}\ \text{sector}: & \text{Complex Scalars}: & (8,1,1,1;1,\overline{4},1,1)+(1,4,1,1;1,1,\overline{4},1)+\\
& & \quad +(1,1,4,1;1,1,1,8)+(1,\overline{4},1,1;8,1,1,1)+\\
& & \quad +(1,1,\overline{4},1;1,4,1,1)+(1,1,1,8;1,1,4,1)\\
& \text{Fermions}_-\ : & \frac{1}{2}(8,1,1,1;8,1,1,1)+(1,4,1,1;1,\overline{4},1,1)+\\
& & \quad +(1,1,4,1;1,1,\overline{4},1)+\frac{1}{2}(1,1,1,8;1,1,1,8)\,.
\end{array}$$

The $\mathbf{28}$ of USp(8) decomposes as $\mathbf{27}+\mathbf{1}$. Again one can check that all gauge and gravitational anomalies cancel.

### 3.2.4 T-duality and the USp(32) string theory

We conclude this section by commenting briefly on an interesting point, concerning the T-duals of the models we have just constructed. Notice that the action of T-duality on anti-D-branes is the same as for D-branes. Therefore, if we perform a T-duality along the four compact directions in the orientifolds of the previous subsections, we obtain a set of models with 32 $\overline{\mathrm{D9}}$-branes and 32 D5-branes, giving rise to the spectra we have presented.

It seems striking to encounter consistent models with 32 $\overline{\mathrm{D9}}$-branes, since in usual type IIB orientifolds, cancellation of untwisted tadpoles imposes the presence of either 32 D9-branes (if $\Omega$ belongs to the orientifold group) or no D9-branes (if it does not). However, reference [24] discussed the existence of an $\Omega$ projection of ten-dimensional type IIB string theory, such that the RR charge of the corresponding crosscaps (that is, the charge of the O9-plane) is opposite to the usual one. Cancellation of tadpoles in this model requires the introduction of $-32$ units of D9-brane charge, which can be minimally accomplished by introducing 32 $\overline{\mathrm{D9}}$-branes. This construction leads to a ten-dimensional tachyon-free non-supersymmetric string theory of open and closed unoriented strings, with USp(32) gauge group. Clearly, the T-duals of the models we have constructed in the previous section correspond to compactifications (with vector structure) of this ten-dimensional non-supersymmetric on $T^4/\mathbb{Z}_N$ spaces. In other words, they correspond to modding out type IIB theory on $T^4/\mathbb{Z}_N$, by the orientifold action $\Omega$ introduced in [24].

The T-duality relation between the type IIB orientifolds of the previous subsections and the type IIB orientifolds with the $\Omega$ projection in [24] seems to suggest interesting relations between the SO(32) and USp(32) theories, at least, after compactification. It would be important to gain a better understanding of such relations. Certainly, the construction of further examples of six-dimensional consistent compactifications of both theories will provide new insights into these issues. We hope our techniques are useful in these investigations.

## 4. Four-dimensional models

In this section we present some examples of four-dimensional type IIB orientifolds containing branes and antibranes. The models are tachyon-free and give rise to chiral spectra. The class of models that can be constructed is presumably very large. The examples we present have been selected to illustrate the following interesting possibility, which has direct application in phenomenological model building. It is possible to construct explicit type IIB orientifolds where $\mathcal{N}=1$ chiral supersymmetric sectors of D-branes are spatially separated (in the internal space) from non-supersymmetric sectors of anti-D-branes. This type of models provides an explicit realization of the supersymmetry breaking scenario where the standard model is embedded in a set of supersymmetric branes, whereas supersymmetry is broken (in our case, at the string scale) in a hidden sector of antibranes. This scenario could be phenomenologically viable for suitable choices of the string and compactification scales, and fits nicely into the circle of ideas recently developed about string theory vacua with large or largish dimensions [17, 18, 19, 20].

### 4.1 A $\mathbb{Z}_6'$ model with non-supersymmetric and supersymmetric sectors

Here we describe a chiral tachyon free four-dimensional example with $\mathcal{N}=1$ supersymmetric sectors, and non-supersymmetric 'hidden' sectors. Branes and antibranes are stuck at different fixed points in the compact space and are unable to move off into the bulk. Therefore, the model is stable against the kind of annihilation and decay mentioned in section 3.1.

The model is based on the $\mathbb{Z}_6'$ orientifold, with twist $v=(1,-3,2)/6$. We will make use of the property, discussed in [9, 29], that tadpoles can be satisfied placing $\mathrm{D5}_3$-branes (and in our case, also $\overline{\mathrm{D5}}_3$-branes) at different fixed points in the second complex plane.

Let us consider the following Chan-Paton matrix for the D9-branes

$$\gamma_{\theta,9} = \mathrm{diag}\left(e^{\pi i \frac{1}{6}}1_4, e^{\pi i \frac{3}{6}}1_8, e^{\pi i \frac{5}{6}}1_4, e^{\pi i \frac{7}{6}}1_4, e^{\pi i \frac{9}{6}}1_8, e^{\pi i \frac{11}{6}}1_4\right). \tag{4.1}$$

Let us also place some D5-branes at the origin in the second complex plane, with Chan-Paton factors

$$\gamma_{\theta,5,0} = \mathrm{diag}\left(e^{\pi i \frac{1}{6}}1_4, e^{\pi i \frac{5}{6}}1_4, e^{\pi i \frac{7}{6}}1_4, e^{\pi i \frac{11}{6}}1_4\right). \tag{4.2}$$

We place $\overline{\mathrm{D5}}$-branes at another fixed point, labeled '1', with matrices

$$\gamma_{\theta,\bar{5},1} = \mathrm{diag}\left(e^{\pi i \frac{3}{6}}1_R, e^{\pi i \frac{9}{6}}1_R\right) \tag{4.3}$$

and some D5-branes at yet another fixed point, labeled '2', with matrices

$$\gamma_{\theta,5,2} = \mathrm{diag}\left(e^{\pi i \frac{3}{6}}1_{8+R}, e^{\pi i \frac{9}{6}}1_{8+R}\right). \tag{4.4}$$

The matrices $\gamma_\Omega$ are defined following the conventions in section 2.1.

The choice of matrices above is just a small variation of the model in [9], and satisfies the tadpole conditions. The spectrum in the 99 sector is the following $\mathcal{N}=1$ theory

$$\begin{gathered}\mathrm{U}(4)\times\mathrm{U}(8)\times\mathrm{U}(4)\\ (6,1,1)+(4,\overline{8},1)+(1,8,\overline{4})+(1,1,6)+(\overline{4},\overline{8},1)+(4,1,\overline{4})+\\ +(1,8,4)+(4,1,4)+(\overline{4},1,\overline{4})+(1,28,1)+(1,\overline{28},1)\,.\end{gathered} \tag{4.5}$$

The D5-branes at the origin form a $\mathcal{N}=1$ supersymmetric sector. For instance, their 55 strings gives the following $\mathcal{N}=1$ supersymmetric spectrum

$$\begin{array}{c} \mathrm{U}(4)\times\mathrm{U}(4) \\ (6,1)+(1,6)+(4,\overline{4})+(4,4)+(\overline{4},\overline{4})\,. \end{array} \tag{4.6}$$

The corresponding 59 matter is given by the $\mathcal{N}=1$ chiral multiplets

$$(4,1;4,1,1)+(1,4;1,\overline{8},1)+(1,\overline{4};1,1,\overline{4})+(\overline{4},1;1,8,1)\,. \tag{4.7}$$

The D5-branes at the fixed point labeled '2' also give an $\mathcal{N}=1$ sector, as follows

|  | $\mathrm{U}(8+R)$ |
|---|---|
| 55 | $\Yasymm+\overline{\Yasymm}$ |
| 59 | $(\Box;\overline{4},1,1)+(\overline{\Box};1,1,4)$ |

(4.8)

Finally, the $\overline{\mathrm{D}5}$-branes give a gauge group $\mathrm{U}(R)$ with the *non-supersymmetric* matter content shown in the following table.

| Sector | Complex scalars | Fermion$_+$ | Fermion$_-$ |
|---|---|---|---|
| $\bar{5}\bar{5}$ | $\Yasymm+\overline{\Yasymm}$ | $\Ysymm+\overline{\Ysymm}$ | Adj. |
| $\bar{5}9+9\bar{5}$ | $(R;4,1,1)+(\overline{R};1,1,\overline{4})$ |  | $(R;\overline{4},1,1)+(\overline{R};1,1,4)$ |

It is easy to check that non-abelian anomalies cancel. Also, U(1) anomalies should cancelled by the GS mechanism proposed in [33].

This model provides a specific realization of a scenario where supersymmetry is broken explicitly at the string scale in an anti-D-brane hidden sector, distant from the $\mathcal{N}=1$ chiral supersymmetric sector. We expect this model to belong to a more general class of explicit type IIB orientifold realizing this type of scenario. Hopefully, this class of models may include theories with more realistic spectra.

In the model we have constructed, if the standard model would be embedded in one of the supersymmetric D5-brane sectors, supersymmetry breaking would be felt through loops involving states in the 99 and $9\bar{5}$ sectors. Therefore, supersymmetry breaking is gauge-mediated. It should not be particularly difficult to construct examples where the supersymmetry breaking is gravity-mediated. This could be accomplished, for instance, by introducing Wilson lines that project all 59 states in the visible sector out.

As mentioned above, branes and antibranes in the model are stuck at fixed points in the compact space, and therefore cannot move off to the bulk and lead to annihilation.[6] However, there exist non-cancelled forces between the branes and antibranes, whose strength depends on the distances between the corresponding fixed points. Another interesting feature of this models is that by increasing $R$ the gauge groups in the model can be made as large as desired, with no apparent inconsistency. The only prize to pay would be an increase in the vacuum energy of the model (cosmological constant) since, roughly speaking, $R$ controls the number of excess brane-antibrane pairs. In our final remarks we speculate on the possible role these two dynamical issues may play in moduli stabilization.

### 4.2 A $\mathbb{Z}_3$ example with non-supersymmetric and supersymmetric sectors

It is easy to construct further examples of four-dimensional models with supersymmetric and non-supersymmetric sectors. Here we discuss a model based on the $\mathbb{Z}_3$ orientifold [14]. In order to cancel the untwisted tadpole the orientifold action requires the introduction of 32 D9-branes and zero net number of D5-branes. However, one can introduce an equal number of D5-branes and $\overline{\text{D5}}$-branes (wrapping, for instance, the third complex plane) and obtain new (non-supersymmetric) consistent models. The tadpole cancellation conditions read

$$\mathrm{Tr}\,\gamma_{\theta,9} + 3(\mathrm{Tr}\,\gamma_{\theta,5,L} - \mathrm{Tr}\,\gamma_{\theta,\bar{5},L}) = -4 \tag{4.9}$$

for any of the nine fixed points in the two first complex planes. Let us consider a particular case, with

$$\gamma_{\theta,9} = \mathrm{diag}\left(1_{16}, e^{2\pi i\frac{1}{3}}1_8, e^{2\pi i\frac{2}{3}}1_8\right), \tag{4.10}$$

that is, $\mathrm{Tr}\,\gamma_{\theta,9} = 8$. This choice implies that no fixed point can be completely empty. Tachyons are most easily avoided by considering the fixed points to contain only D5-branes or only $\overline{\text{D5}}$-branes. We label such fixed points by indices $M$, $K$, respectively. Twisted tadpoles (4.9) require $\mathrm{Tr}\,\gamma_{\theta,5,M} = -4$, $\mathrm{Tr}\,\gamma_{\theta,\bar{5},K} = 4$, therefore

$$\begin{aligned}
\gamma_{\theta,5,M} &= \mathrm{diag}\left(1_{n_M}, e^{2\pi i\frac{1}{3}}1_{n_M+4}, e^{2\pi i\frac{2}{3}}1_{n_M+4}\right) \\
\gamma_{\theta,\bar{5},K} &= \mathrm{diag}\left(1_{m_K+4}, e^{2\pi i\frac{1}{3}}1_{m_K}, e^{2\pi i\frac{2}{3}}1_{m_K}\right).
\end{aligned} \tag{4.11}$$

Since no net D5-brane charge is allowed

$$\sum_M (3n_M + 8) = \sum_K (3m_K + 4)\,. \tag{4.12}$$

[6] Upon a closer look, a possible process of anihilation may take place in the region of moduli space where the points labeled '1' and '2' are very close, and a winding mode stretched between branes and antibranes becomes tachyonic. Even though the anihilation would involve objects at different fixed points, it seems to be consistent with tadpole cancellation in the final state. This new kind of instability is not generic and is absent in other models, in particular the $\mathbb{Z}_3$ orientifold in the next section.

A particular solution is obtained by choosing $n_M = 0$, $m_K = 0$, and three fixed points containing only D5-branes, and six with only $\overline{\rm D5}$-branes. This choice has the virtue that all branes and antibranes are stuck at the fixed points.

The spectrum in the general case can be computed using our rules. Since it is not particularly interesting, we spare the reader these details. Let us just mention that the gauge group of the model is

$$\mathrm{SO}(16) \times \mathrm{U}(8) \times \prod_M [\mathrm{USp}(n_M) \times \mathrm{U}(n_M+4)] \times \prod_K [\mathrm{USp}(m_K+4) \times \mathrm{U}(m_K)] \quad (4.13)$$

and the spectrum of fermions is

$$\begin{array}{lll}
\mathbf{99} & \mathrm{Fermions}_+ & 3(16,\overline{8}) + 3(1,28) \\
 & \mathrm{Fermions}_- & \left(\begin{smallmatrix}\square\\\square\end{smallmatrix},1\right) + (1,\mathrm{Adj}) \\
\mathbf{5_M 9} & \mathrm{Fermions}_+ & (n_M,1;1,\overline{8}) + (1,n_M+4;1,8) + (1,\overline{n_M+4};16,1) \\
\mathbf{\bar{5}_K 9} & \mathrm{Fermions}_- & (m_K+4,1;1,\overline{8}) + (1,m_K;1,8) + (1,\overline{m_K};16,1) \\
\mathbf{5_M 5_M} & \mathrm{Fermions}_+ & 3(n_M,\overline{n_M+4}) + 2\left(1,\begin{smallmatrix}\square\\\square\end{smallmatrix}\right) + (1,\square\!\square) \\
 & \mathrm{Fermions}_- & (\square\!\square,1) + (1,\mathrm{Adj}) \\
\mathbf{\bar{5}_K \bar{5}_K} & \mathrm{Fermions}_+ & 3(m_K+4,\overline{m_K}) + 2(1,\square\!\square) + \left(1,\begin{smallmatrix}\square\\\square\end{smallmatrix}\right) \\
 & \mathrm{Fermion}_- & \left(\begin{smallmatrix}\square\\\square\end{smallmatrix},1\right) + (1,\mathrm{Adj})\,.
\end{array} \quad (4.14)$$

One can check that all anomalies cancel (to show that, one has to use the fact that none of the nine fixed points in the first two complex planes is empty).

Therefore we obtain a set of models with diverse numbers of supersymmetric and non-supersymmetric sectors. Most of the comments in the previous section apply to these models as well, most notably the gauge-mediated nature of supersymmetry breaking in the supersymmetric D5-brane sectors.

We hope the examples in this section suffice to illustrate the construction of consistent four-dimensional orientifolds with antibranes, and the relative ease with which supersymmetric and non-supersymmetric sectors within the same construction can be implemented

## 5. Conclusions

In this paper we have provided the basic rules for the construction of type IIB orientifold models with branes and antibranes. The models present the attractive feature that closed string tachyons are automatically absent, since supersymmetry is preserved in the closed string sector. Moreover, open string tachyons from strings stretching between branes and antibranes can be avoided. Therefore many new tachyon-free non-supersymmetric models can be constructed using this simple idea. We have provided several explicit examples in six and four dimensions, which illustrate the generic features of these models.

A particular interesting application of these constructions, from the phenomenological viewpoint, is that they allow the presence of $\mathcal{N}=1$ supersymmetric sectors of branes and supersymmetry breaking sectors of antibranes, spatially separated in the compact space.

The models may also play an interesting role in the web of string dualities involving other non-supersymmetric strings. We have discussed that type IIB orientifolds on $T^4/\mathbb{Z}_N$ with vector structure are related by T-dualtity to analogous compactifications of the non-supersymmetric USp(32) theory of [24]. We hope further research on these vacua will uncover new interesting properties.

Finally, we would like to point out two important issues we have not addressed, concerning the dynamics of moduli of these models. As briefly mentioned, branes and antibranes in the model suffer attractive forces due to the lack of supersymmetry. If these objects are stuck at different fixed points, their attraction presumably leads to a potential for the compactification moduli, pushing their vevs to small radii. If the models contain other sources of non-trivial moduli dynamics, pushing the radii to large values, this effect would contribute to their stabilization. The second point refers to the cosmological constant. Computing the vacuum energy by simply taking into account the tension of branes, antibranes and orientifold planes (*i.e.* ignoring for the moment interaction energies), it is easy to see that the models we have discussed have a positive cosmological constant, which is, roughly, inversely proportional to the string coupling constant. Even without having addressed its specific value, which furthermore would be rather model-dependent, it is tempting to interpret this vacuum energy as a potential for the dilaton. From that point of view, the model would force the string coupling to become strong. If the model has sources of dilaton potential with a runaway behaviour to small coupling (for instance, a gaugino condensate), the dilaton would be stabilized.

In any case, interesting dynamics for the moduli seems to arise from the absence of supersymmetry. Whether stabilization takes place, and whether the resulting values for the compactification scale and string coupling are of phenomenological interest, depends on dynamical information we have not yet studied. We hope further reseach in these models succeeds in uncovering their rich dynamics and their phenomenological potential.

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