Discrete Deformations in Type I Vacua

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Abstract

We study supersymmetric orientifolds where the world-sheet parity transformation is combined with a conjugation of some compact complex coordinates. We investigate their T-duality relation to standard orientifolds and discuss the origin of continuous and discrete moduli. In contrast to standard orientifolds, the antisymmetric tensor describes a continuous deformation, while the off-diagonal part of the metric is frozen to quantized values and is responsible for the rank reduction of the gauge group. We also give a geometrical interpretation of some recently constructed six-dimensional permutational orientifolds.
1. Introduction

Since the early work of Bianchi, Pradisi and Sagnotti \[1\], it is known that Type I string vacua involve both continuous and discrete moduli. Thus, in the case of ordinary orientifolds based on the world-sheet parity $\Omega$ \[2\], discrete internal backgrounds for the NS-NS two-form, $B_{ab}$, are still allowed, although the corresponding (continuous) deformations are projected out. A background $B$-field generically leads to a rank reduction for the gauge group, while, in the case of orbifold compactifications, extra multiplicities appear in some open string sectors \[3\].

In a recent paper \[4\] a new class of six-dimensional orientifolds of the type IIB string has been studied. They involve the combined action of the world-sheet parity $\Omega$ with a conjugation $\mathcal{R}$ of the internal complex coordinates and a geometric $\mathbb{Z}_N$ action. The new feature arising in these orientifolds is that the Klein bottle amplitude receives contributions from all $\mathbb{Z}_N$ twisted sectors. After an $S$ transformation to the transverse tree-channel, these lead to new tadpoles whose cancellation requires (different types of) D7-branes intersecting at non-trivial (quantized) angles. Then, strings stretched between different D7-branes give rise to new twisted open sectors crucially involved in tadpole cancellations. The resulting models have $\mathcal{N} = (1, 0)$ space-time supersymmetry in $D = 6$ and are characterized by Chan-Paton groups of reduced rank. While for the standard $\Omega$ projection this phenomenon may be related to the presence of a quantized background for the NS-NS antisymmetric tensor (both in toroidal \[1\] and in orbifold \[3\] compactifications), a similar understanding is missing for the $\Omega \mathcal{R}$ orientifolds of \[4\]. Finding an appropriate description of the observed rank reduction is the main motivation of the present letter.

As a simple example, in Section 2 we study $\Omega \mathcal{R}$ orientifolds of type IIA on a two-torus and using T-duality we gain an understanding of the mechanism responsible for the rank reduction. In Section 3 we study tadpole cancellation for this model, derive the quantization rules for geometric moduli and show that they are consistent with the results obtained in \[4\]. Finally, in Section 4 we study the six dimensional $\mathbb{Z}_2 \Omega \mathcal{R}$ orientifold and using T-duality give an alternative geometric interpretation of the extra multiplicities present in the Neuman-Dirichlet open-string sector.
2. T-duality and $\Omega R$ orientifolds

Let us consider the type IIB string on a torus $T^2$ with complex coordinate $z = x_2 + ix_1$. Without loss of generality, one can choose the basis

\[ e'_1 = iR, \quad e'_2 = -b + ia, \]  

(2.1)

in terms of which the complex and Kähler structures are

\[ U' = U'_1 + iU'_2 = \frac{e'_2}{e'_1} = \frac{a}{R} + i\frac{b}{R}, \quad T' = T'_1 + iT'_2 = B' + iV' = B' + ibR, \]  

(2.2)

where we have also turned on an internal background $B$-field. It is well known that modding out by the world-sheet parity transformation $\Omega$, the three geometric moduli $U'_1, U'_2, T'_2$ survive the projection whereas, up to the identification $B' \equiv B' + \mathbb{Z}$, the antisymmetric tensor is frozen to the two possible values $B' = 0$ and $B' = \frac{1}{2}$ [1].

A T-duality along the $x_1$ direction leads to the orientifold

\[ \text{type IIA on } T^2 \left/ \Omega R \right., \]  

(2.3)

where now the $\Omega$ projection is combined with the operation $\mathcal{R} : z \rightarrow \bar{z}$. Moreover, the complex and Kähler structures are interchanged, so that for the dual torus $T^2$

\[ U = U_1 + iU_2 = \frac{e_2}{e_1} = \frac{-b + iB'/R}{i/R}, \quad T = T_1 + iT_2 = B + iV = \frac{a}{R} + i\frac{b}{R}. \]  

(2.4)

Therefore, in the orientifold (2.3) the $B$-field is a continuous parameter while the geometric modulus $U_1$ is frozen. The quantization condition $2U_1 \in \mathbb{Z}$ expresses the fact that for all other values of $U_1$, the action of $\mathcal{R}$ on $T^2$ is not crystallographic. Moreover, under T-duality the periodicity $B' \equiv B' + \mathbb{Z}$ is mapped to the invariance of the complex structure under integer $\text{SL}(2; \mathbb{Z})$ shifts.

In conclusion, from T-duality we learn that for $2U_1 \in 2\mathbb{Z} + 1$ the rank of the gauge group is halved. In the next Section, for the orientifold (2.3) we discuss tadpole cancellation in detail and confirm the above expectations suggested by T-duality.
3. Tadpole cancellation for $\Omega R$ orientifolds

Let us therefore consider the compactification on a generic two-torus $T^2$, whose complex ($U = U_1 + i U_2$) and Kähler ($T = T_1 + i T_2$) structures are related to the metric and antisymmetric tensor by

$$g = \frac{\alpha' T_2}{U_2} \begin{pmatrix} 1 & U_1 \\ U_1 & |U|^2 \end{pmatrix}, \quad B = \alpha' T_1 \begin{pmatrix} 0 & -1 \\ 1 & 0 \end{pmatrix}.$$  \hspace{1cm} (3.1)

The left and right momenta corresponding to the complex coordinates ($z, \bar{z}$) on the $T^2$ are

$$p_L = \frac{1}{\sqrt{\alpha' U_2 T_2}} \left[ U m_1 - m_2 - \bar{T} (n^1 + U n^2) \right], \quad p_R = \frac{1}{\sqrt{\alpha' U_2 T_2}} \left[ U m_1 - m_2 - T (n^1 + U n^2) \right],$$  \hspace{1cm} (3.2)

where $m_a$ and $n^b$ denote the KK momenta and winding numbers. For generic values of the $U$ and $T$ moduli, the type IIB string is no more left-right symmetric. However, one can impose the invariance under $\Omega$ and find constraints on the allowed values of the moduli. In particular, in the “conventional” type I string, obtained projecting the parent IIB with respect to the simple world-sheet parity [2], the real slice of the Kähler structure (the $B$-field) can only take discrete values ($2 T_1 \in \mathbb{Z}$) [1]. This is consistent with the fact that the excitations of the NS-NS two-form are no longer part of the physical spectrum of the projected theory. Moreover, the presence of a quantized flux for the $B$-field implies the reduction of the total size of the Chan-Paton gauge group by a factor $2^{r/2}$, with $r = \text{rank}(B_{ab})$, and a continuous interpolation (via Wilson lines) between orthogonal and symplectic gauge groups [1].

Actually, one can consider more general projections and dress the world-sheet parity $\Omega$ with other type II symmetries as, for instance, coordinate reflections (T-dualities) in the compact directions. For the two-dimensional torus there exist two additional possibilities, namely the inversion of both coordinates, $I: z \to -z$, or of a single coordinate

$$\mathcal{R}: z \to \bar{z}.$$  \hspace{1cm} (3.3)

It is not difficult to see that $\Omega I$ does not present any substantial difference compared to the standard $\Omega$ projection. In both cases, only the imaginary slice of the Kähler structure survives, while the resulting type I models correspond to compactifications on T-dual tori (with D7 and D9-branes, respectively).
However, new interesting features arise when the world-sheet parity is dressed with the conjugation \((3.3)\). A simple analysis of the massless spectrum reveals that the internal components of the NS-NS antisymmetric tensor

\[
\left(\psi_{\frac{1}{2}} - \psi_{-\frac{1}{2}} \tilde{\psi}_{\frac{1}{2}} - \psi_{\frac{1}{2}} \tilde{\psi}_{-\frac{1}{2}}\right) |00\rangle \quad (3.4)
\]

now survive the \(\Omega R\) projection, while the mixed components of the internal metric

\[
\left(\psi_{\frac{1}{2}} \tilde{\psi}_{-\frac{1}{2}} + \psi_{-\frac{1}{2}} \tilde{\psi}_{\frac{1}{2}}\right) |00\rangle \quad (3.5)
\]
do not. Therefore, in this case one expects that the antisymmetric tensor is a continuous modulus of the projected theory, while some conditions have to be satisfied by the mixed components of the metric. Indeed, it is not hard to see that requiring invariance of the parent theory under

\[
\Omega R : \quad p_L = \tilde{p}_R \quad (3.6)
\]
results in a quantization condition for the real slice of the complex structure

\[
2U_1 \in \mathbb{Z} \quad \text{,} \quad (3.7)
\]
in analogy with the standard case.

Once the \(\Omega R\) symmetry of the parent closed string is restored, one can proceed as in [2] to construct the open descendants. Starting as usual from the Klein bottle amplitude

\[
\mathcal{K} = \frac{1}{2} \text{Tr} \left(\Omega R \mathcal{P}_{GSO} e^{-2\pi t(L_0 + \bar{L}_0)}\right) \quad (3.8)
\]
one realizes that it receives contributions from the lattice states satisfying \(p_L = \tilde{p}_R\), which are fixed under \(\Omega R\). The resulting amplitude in the direct loop-channel is

\[
\mathcal{K} = \frac{1}{2} (V_S - S_S)(2it) \sum_{m_2,n_1} \frac{e^{-\frac{\pi t}{T_2} |m_2 + T n_1|^2}}{\eta^2(2it)} , \quad (3.9)
\]
and after an \(S\) transformation to the transverse channel leads to the amplitude

\[
\tilde{\mathcal{K}} = \frac{\eta^5}{2} U_2 (V_S - S_S)(i\ell) \sum_{m_1,n_2} \frac{e^{-2\pi \ell \frac{T_2}{T_1} |m_1 - T n_2|^2}}{\eta^2(i\ell)} . \quad (3.10)
\]

In \((3.11)\) we have written the contribution of the world-sheet fermions

\[
(1 - 1) \frac{\eta^4}{2\eta^4} = V_S - S_S \quad (3.11)
\]
in terms of level-one SO(8) characters, and we have omitted the contributions of the
transverse bosonic coordinates and the integration measure.

The annulus amplitude can be computed in two ways. The first, as in [4], is to place
D8-branes at \( x_1 = 0 \) extending in the \( x_2 \) direction and compute the direct-channel amplitude
taking care exactly of the KK and winding states in the \( x_2 \) and \( x_1 \) directions. The
other, followed here, is to start from the transverse channel, that involves only the sectors
of the parent theory that are combined with their conjugates. For the \( \Omega R \) projection,
these states are selected imposing \( p_L = -\bar{p}_R \). Given the quantization condition (3.7) on
the \( U_1 \) field, there are further restrictions on these allowed states:

\[
2U_1 m_1, \ 2U_1 n^2 \in \mathbb{Z}.
\]

As usual, these constraints are imposed inserting in the transverse tree-channel annulus
amplitude suitable projectors

\[
\tilde{A} = \frac{2^{5+2a-2}}{2} U_2 N^2 (V_8 - S_8)(i\ell) \sum_{\epsilon_1,\epsilon_2 = 0,1} \sum_{m_1, n^2} e^{-\frac{\pi U_1}{U_2} m_1 - T n^2} e^{2i\pi U_1 (m_1 \epsilon_1 - n^2 \epsilon_2)} \eta^2(i\ell),
\]

where \( a \) depends on the value of \( U_1 \), and is zero for \( U_1 \in \mathbb{Z} \), and one for \( U_1 \in \mathbb{Z} + \frac{1}{2} \).
The coefficient \( 2^{2a-2} \) ensures a proper normalization of the vector in the direct-channel
amplitude

\[
A = \frac{2^{2a-2}}{2} N^2 (V_8 - S_8)(it/2) \sum_{\epsilon_1,\epsilon_2 = 0,1} \sum_{m_2, n} e^{-\frac{\pi U_1}{U_2} m_2 + U_1 \epsilon_2 + T (n^2 + U_1 \epsilon_2)} \eta^2(it/2),
\]

that, for \( U_1 \in \mathbb{Z} + \frac{1}{2} \), receives contribution only from \( \epsilon_1 = \epsilon_2 = 0 \). The M"obius amplitude

\[
\tilde{M} = -\frac{2 \times 2^{a-1}}{2} U_2 N (\hat{V}_8 - \hat{S}_8)(i\ell + \frac{1}{2}) \times
\sum_{\epsilon_1,\epsilon_2 = 0,1} \sum_{m_1, n^2} e^{-2\pi \ell \epsilon_2 U_1} e^{2i\pi U_1 (m_1 \epsilon_1 - n^2 \epsilon_2)} \gamma_{\epsilon_1,\epsilon_2} \eta^2(i\ell + \frac{1}{2})
\]

involves signs \( \gamma_{\epsilon_1,\epsilon_2} \) as in [4], that enforce a proper normalization of the \( U_1 \) projector.

Extracting the contributions of \( \tilde{K}, \ \tilde{A} \) and \( \tilde{M} \) to the massless tadpoles, one finds the
consistency condition

\[
\sum_{\epsilon_1,\epsilon_2 = 0,1} \gamma_{\epsilon_1,\epsilon_2} = 2,
\]

\[5\]
and, as a result, the total size of the Chan-Paton gauge group is reduced by a factor $2^a$

$$N = 2^{5-a}.$$  \hspace{1cm} (3.17)

Thus, the gauge group is SO($2^{5-a}$) or USp($2^{5-a}$) depending on the sign of $\gamma_{0,0}$ in

$$\mathcal{M} = -\frac{2^{a-1}}{2}N (\hat{V}_8 - \hat{S}_8) ((1 + it)/2) \sum_{\epsilon_1, \epsilon_2 = 0,1} \sum_{m_2, n_1} e^{-\frac{\pi i}{v_2 T_2} |m_2 + U_1 \epsilon_2 + \bar{T}(n_1 + U_1 \epsilon_1)|^2} \gamma_{\epsilon_1, \epsilon_2}. \hspace{1cm} (3.18)$$

For $U_1 \in \mathbb{Z}$ its rank is not reduced, while for $U_1 \in \mathbb{Z} + \frac{1}{2}$ it is halved. Also in this case Wilson lines connect continuously orthogonal and symplectic groups passing through unitary ones.

We can now turn to analyze more complicated cases. In [4], six-dimensional $\Omega R$ orientifolds with $\mathcal{N} = (1,0)$ space-time supersymmetry were studied. In particular, the orientifolds

$$\text{type IIB on $T^2 \times T^2$} \quad \frac{\{\Omega R, \mathbb{Z}_N\}}{}$$  \hspace{1cm} (3.19)

with $N = 3, 4, 6$ were considered. In order to define perturbatively consistent orientifolds, the lattices had to allow a crystallographic action of the $\mathbb{Z}_N$ orbifold group and had also to be oriented in specific ways relative to the mirror plane of $R$. The massless spectra of the orientifolds (3.19) had the generic feature of rank reductions of the gauge groups by powers of two. For instance, for the $\mathbb{Z}_3$ model a natural choice was the lattice $\mathbf{A}$ depicted in Fig. 1, both in the (67) and (89) planes.

Fig. 1. The $T^2$ lattices $\mathbf{A}$ for the $\mathbb{Z}_3$ orbifold.
Now, after applying an SL(2; \mathbb{Z}) transformation on each $T^2$ leading to a different unit cell with base vectors
\[ e_1 = iR, \quad e_2 = -\frac{\sqrt{3}}{2} + i\frac{R}{2}, \] (3.20)
one can see that the background involves a quantized value $U_1 \in \mathbb{Z} + \frac{1}{2}$ of the real component of the complex structure. One thus expects a rank reduction of the gauge group by a factor of two for each $T^2$, indeed consistent with the SO(8) gauge group derived in [4] for this model. For the tadpole cancellation condition and the computation of the massless spectrum, it is important to know how the $\mathbb{Z}_3$ fixed points transform under the reflection $\mathcal{R}$. For the lattice in Fig. 1 one of the nine fixed points is invariant while the remaining eight are interchanged.

![Fig. 2. The $T^2$ lattices B for the $\mathbb{Z}_3$ orbifold.](image)

Rotating the lattice by an angle $\phi = \pi/6$ results in the unit cell $B$ of Fig. 2. By an SL(2; \mathbb{Z}) transformation the unit cell can be brought to the form
\[ f_1 = iR, \quad f_2 = -\frac{1}{2\sqrt{3}} + i\frac{R}{2}. \] (3.21)

Note, that for the lattice $B$ all three $\mathbb{Z}_3$ fixed points of a single $T^2$ are invariant under $\mathcal{R}$. Thus, besides the $T^4$ torus $AA$ there exist two more inequivalent choices, namely $AB$ and $BB$. Going through the computation eventually leads to the anomaly-free massless spectra shown in Table 1. Particularly interesting is the $BB$ model since it corresponds to a dual heterotic string with a frozen dilaton. Notice, that all these spectra ($AA$, $AB$, $BB$) have already appeared in the study of open descendants of six-dimensional Gepner models.
Moreover, the AB case was recently obtained in a study of asymmetric orientifolds [7]. In the Gepner construction, the various choices of elementary cells correspond to different modular invariant combinations of characters for the parent closed theory. A better understanding of this correspondence could shed some light on the D-brane interpretation of orientifolds of Gepner models [8,9].

| $T^4$ | closed | open |
|-------|--------|------|
| AA    | $13H + 8T$ | $SO(8) + 2H$ in 28 |
| AB    | $15H + 6T$ | $SO(8) + 4H$ in 28 |
| BB    | $21H$ | $SO(8) + 10H$ in 28 |

Table 1: String spectra of $T^4/Z_3$

For the $Z_4$ orientifold, one had to make the choices in Figs. 3 and 4 for the lattices in the (89) and (67) directions, that correspond to the unit cells

$$
e_1 = iR_1, \quad e_2 = -R_1, \quad f_1 = iR_2, \quad f_2 = -\frac{R_2}{2} + i\frac{R_2}{2}.
$$

(3.22)

Fig. 3. The $Z_4$ lattice A in the plane (89).

In analogy with the previous case, one would thus expect that the rank of the gauge group is halved for this model. Since $Z_4$ contains a $Z_2$ subgroup, tadpole cancellation requires the introduction of two unrelated types of D7-branes and the unreduced rank of the gauge group is 32. Consistently, the $Z_4$ orientifold was found to have the gauge group $U(8) \times U(8)$ of rank 16.
Finally, for the $\mathbb{Z}_6$ orientifold the only consistent choice of the lattices was the $\mathbb{Z}_3$ AB configuration leading to a reduction of the gauge group by a factor of four. This is nicely confirmed by the resulting gauge group $U(4) \times U(4)$ found in [4].

Four dimensional orientifolds with $\mathcal{N} = 1$ supersymmetry discussed in the forthcoming paper [10] display gauge groups that precisely reflect the quantization rules for $U_1$.

4. The $\mathbb{Z}_2$ $\Omega R$ orientifolds

As we anticipated in Section 2, toroidal $\Omega$ orientifolds are related by T-duality to toroidal $\Omega R$ ones, where Kähler and complex structures are interchanged. However, for orientifolds of toroidal orbifolds the situation is a bit more involved, as under T-duality, $(z_L, z_R) \rightarrow (\bar{z}_L, z_R)$, in general the left-right symmetric $\mathbb{Z}_N$ action is mapped to a left-right asymmetric $\hat{\mathbb{Z}}_N$

$$(z_L, z_R) \rightarrow (e^{\frac{2\pi i}{N}} z_L, e^{-\frac{2\pi i}{N}} z_R).$$

(4.1)

Thus, in general T-duality relates symmetric with asymmetric orientifolds, but for the $\mathbb{Z}_2$ orbifold the action (4.1) remains left-right symmetric. In this particular case T-duality identifies

$$\begin{align*}
\text{type IIB on } T^2 & \times T^2 \\
\{\Omega, \mathbb{Z}_2\} & \quad (4.2)
\end{align*}$$

with

$$\begin{align*}
\text{type IIB on } T^2 & \times T^2 \\
\{\Omega R, \mathbb{Z}_2\} & \quad (4.3)
\end{align*}$$

As a result, the two independent values of the quantized $B$-field for the $T^2$ factors in (4.2) are mapped to the corresponding $A$ and $B$ tori in Figs. 3 and 4. For instance, the $\Omega$ orientifold with vanishing $B$-field corresponds to the $\Omega R$ orientifold on the $AA$ torus, as indeed can be confirmed by a direct computation.
Turning on a $B$-field, or equivalently rotating the cells, leads to the expected rank reduction of the gauge group and introduces multiplicities in the Neuman-Dirichlet sector, that are crucial for the cancellation of the anomalies by a generalized Green-Schwarz mechanism \cite{11}. In \cite{3}, these multiplicities had been related to the (modified) structure of the fixed points in the presence of a non-vanishing $B$-flux, and indeed the fixed points of the $\mathbb{Z}_2$ orbifold fill multiplets of dimension $2^{r/2}$. Thus, a D5-brane in the presence of a rank $r$ $B_{ab}$ corresponds to $2^{r/2}$ (equivalent) copies, each sitting at a fixed point in the multiplet, while the multiplicity $2^{r/2}$ in the 95 sector reflects the number of equivalent ways of building an ND string.

An alternative geometric description can now be given in the context of $\Omega\mathcal{R}$ orientifolds. Under T-duality, the D9 and D5 branes are mapped to D7-branes stretching along the horizontal and vertical directions in Figs. 3 and 4, respectively. In the $A$ torus the intersection number of these two D7-branes is one while it is two for the $B$ torus$^1$. Thus, in the $\Omega\mathcal{R}$ orientifolds, the extra multiplicities in the 95 sector translate into the intersection numbers of the corresponding D7-branes: one for the $AA$ lattice (corresponding to $r = 0$), two for the $AB$ lattice (corresponding to $r = 2$) and finally four for the $BB$ lattice (corresponding to $r = 4$).

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$^1$ The relation between these intersection numbers and the extra open string multiplicities $\kappa_k$ introduced in \cite{4} is explained in more detail in \cite{10}.
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