We discuss the interplay between freely acting orbifold actions, discrete deformations and internal uniform magnetic fields in four-dimensional orientifold models.

1 Introduction

In the framework of string/M-theory attempts to single-out models with a particle content as close as possible to (some extension of) the Standard Model, a prominent role is played by type I vacua with internal magnetic fields or, in a T-dual language, with D-branes intersecting at angles. In this talk, we review the basic ingredients entering these constructions, stressing their most appealing features, namely chirality, the natural emergence of replicas of matter-field families and the surprising power of combinations of shifts in freely acting orbifolds and the NS-NS $B_{ab}$, itself equivalent to an asymmetric shift-orbifold projection. In particular, after a general discussion, we show an example that displays all these properties, extracted from ref. (8), where an exhaustive discussion of four dimensional $Z_2 \times Z_2$ shift-orientifolds can be found.

2 Shift-Orientifolds

The simplest example of a shift orbifold can be obtained starting from a one-dimensional boson $X$ compactified on a circle of radius $R$ and identifying its values under the (freely acting) antipodal transformation, namely $X \rightarrow X + \pi R$. Twisted sectors depend on $R$, and the orbifold procedure, not surprisingly, yields simply a one-dimensional boson compactified on a circle of radius $R/2$. However, the combination of freely-acting actions with internal (discrete) symmetries realizes in string theory the Scherk-Schwarz mechanism via the assignment of different periodicity conditions to bosons.

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and fermions along the internal directions, giving rise to the (spontaneous) breaking of supersymmetry. In the framework of type II superstrings, a T-duality can turn a shift of internal momenta into a shift of internal windings, that therefore exhibit analogous features. In orientifolds, however, things are quite different. Indeed, momentum shifts are conventional Scherk-Schwarz deformations, while winding shifts can be related via T-duality to momentum shifts along the 11-th coordinate of the underlying M-theory. Alternatively, the “massless” spectrum depends crucially on the relative orientations of the momentum shifts and the D-branes present. Thus, a shift longitudinal to the D-branes in a certain sector gives rise to a conventional Scherk-Schwarz deformation, with supersymmetry already broken at tree level in string perturbation theory. On the contrary, with a momentum shift orthogonal to the D-branes in a certain sector, the supersymmetry of the massless modes (and sometimes also of the massive ones) is not affected by the orbifold action. This phenomenon, termed “brane supersymmetry” in the literature, has a very neat geometric interpretation: it can be ascribed to configurations with multiplets of D-branes. Indeed, an orthogonal momentum shift is compatible with a brane configuration if, and only if, for each given D-brane at a certain “position” in spacetime, there exist image D-branes at each image “position” obtained from the initial one. Multiplets of branes are a generic feature of shift-orientifolds. A very interesting class of these models in four dimensions, extensively studied in refs. (11) and (12), corresponds to orientifolds of freely acting orbifolds obtained combining $Z_2 \times Z_2$ involutions and shifts. More precisely, the six torus is taken as a product of three two-tori $T^2 \times T^2 \times T^2$, while the orbifold action is a combination of the $Z_2 \times Z_2$ group containing the identity and the three involutions of pairs of the three complex torus coordinates, combined with shifts along their real parts. Up to T-duality and trivial re-

| models      | $D_9$ susy | $D_5_1$ susy | $D_5_2$ susy |
|-------------|-----------|-------------|-------------|
| $p_3$       | N=1       | N=2         | N=2         |
| $w_2p_3$    | N=2       | N=2         | N=4         |
| $w_1w_2p_3$ | N=4       | N=4         | N=4         |
| $p_2p_3$    | N=1       | N=2         | -           |
| $w_1p_2$    | N=2       | N=4         | -           |
| $w_1^3p_3$  | N=2       | N=4         | -           |
| $w_1^2p_3w_3$| N=4     | N=4         | -           |
| $p_1p_2p_3$ | N=1       | -           | -           |
| $p_1w_2w_3$ | N=2       | -           | -           |
| $w_1w_2w_3$ | N=4       | -           | -           |
definitions, the types of inequivalent models are collected in table 1, together with their degree of “brane supersymmetry”, that signals the brane-multiplet structure of the corresponding orientifolds. Their open spectra are very rich but unfortunately are not chiral, and can be found in refs. (11) and (12).

3 B_{ab} and Shifts

The presence of a NS-NS two form field B_{ab} in type I vacua has far reaching consequences that induce several interesting effects. The orientifold projection is actually compatible only with quantized values of the background field, a feature that can be easily deduced analyzing the type IIB superstring compactified on a d-dimensional torus. The generalized left and right momenta entering the Narain lattice can be written in the form

\begin{align}
    p_{L,a} &= m_a + \frac{1}{\alpha'} (g_{ab} - B_{ab}) n^b, \\
    p_{R,a} &= m_a - \frac{1}{\alpha'} (g_{ab} + B_{ab}) n^b, \quad (1)
\end{align}

where m_a and n^a are the (integer) momentum and winding quantum numbers, respectively. The Narain lattice admits a sensible action of the world-sheet parity operator Ω only if the condition

\[ \frac{2}{\alpha'} B_{ab} \in \mathbb{Z} \]

holds\(^a\): B_{ab} must belong to the integer cohomology of the d-dimensional torus and the independent values of its components are 0 and 1/2 in suitable inverse α' units. Equivalently, the quantized values of B_{ab} can be ascribed to a toroidal compactification without vector structure. Indeed, the deep connection between the bulk (closed) sector and the boundary (open) sector in orientifolds reflects itself in the identification of the quantized B_{ab}, a theta-like angle of the closed-string sector, with the mod-two cohomology class \( \tilde{w}_2 \) that, in analogy with the Stiefel-Whitney class \( w_2 \) for spin-bundles, measures the obstruction to defining a vector structure on the vacuum gauge bundle. Gauge fields are in the brane sectors of orientifolds, but their numbers and representations depend on the presence of non-trivial B_{ab} fluxes. The most striking effect of a non-vanishing B_{ab}-field on orientifolds is the reduction of the rank of the Chan-Paton gauge group by a factor of \[ 2^r/2 \], where r is the rank of B_{ab}. In the original derivation\(^13\), the rank reduction was deduced

\(^a\)If the orientifold projection is a combination of Ω and an involution, the discretization affects other moduli, for instance some off-diagonal components of the torus metric.\(^14\)
algebraically considering the one loop partition function at rational points in the moduli space, where the quantized $B_{ab}$ was implicit in the Narain lattice choice\cite{footnote3} and at generic radii, imposing the correct normalization of the closed states along the tube in the annulus and Möbius amplitudes. The topological properties of the toroidal compactification without vector structure allow an alternative interpretation in terms of non-commuting Wilson lines. These feel the $\tilde{w}_2$, leading to a non-commutativity of the Chan-Paton gauge matrices, that reduces by the mentioned factor $2^{r/2}$ the rank of the Chan-Paton group\cite{footnote15,footnote16,footnote17,footnote18,footnote19,footnote20}. $r$ is also the parameter that discriminates between disconnected components of the moduli space of type I compactifications with 16 supercharges\cite{footnote20}.

Geometric interpretations can be given in terms of the extended objects present in the target space of type I vacua, namely $D$-branes and $O$-planes. Without loss of generality, we can limit ourselves to the case of a two-torus equipped with space-filling $O_{2\pi}$-planes and stacks of $N$ $D_2$ branes, that correspond to configurations of standard toroidal compactifications \textit{with} vector structure and a standard $\Omega$ orientifold projection. More complicated cases exist, but the qualitative features remain the same. It was observed in ref.\cite{footnote19} that, performing a pair of T-dualities along the two directions of the torus,

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig1}
\caption{$O_0$-planes configuration on a two-torus: a) $B = 0$. b) $B \neq 0$}
\end{figure}
one ends up with a configuration exhibiting a quartet of $O_{0^+}$-planes located at the fixed points of a $Z_2$ orbifold of the two-torus itself, so that, say, $N$ $D$-branes are required to neutralize the overall R-R charge of the configuration. On the other hand, if the two-torus is without vector structure, two different types of $O$-planes are present, the conventional $O^+$’s with tension and charge opposite to those of the $D$-branes, and “exotic” $O^-$-planes with reverted tension and charge. As a result, the net charge of the configuration is reduced by a factor of two, because after two T-dualities one is faced with three $O^+$-planes and one $O^-$-plane (see fig. 1), in a configuration that requires a net number of $N/2$ $D$-branes (as compared to $N$, the number previously required) to neutralize the total R-R charge. All this reflects precisely the features of partition functions in ref. [13].

An interesting alternative geometrical interpretation, useful in the context of intersecting $D$-brane models, can also be given as follows. Let us start, for clarity, with a rectangular torus of sides $R_1$ and $R_2$, and let us perform a T-duality along the vertical direction in such a way that the complex and Kähler structures get interchanged. The result depends crucially on the vector structure of the torus. Indeed, as shown in fig. 2, the $O^2$’s and the $D^2$’s becomes $O^1$’s and $D^1$’s after the T-duality, but the shape of the final torus is very different in the two cases: if $B_{ab} = 0$, one is left with a new rectangular torus of sides $R_1$ and $R_2' = \alpha'/R_2$. On the contrary, if $B_{ab} \neq 0$, one ends up with a tilted torus admitting vector structure, i.e. with $B_{ab}' = 0$. However, in the latter case, each of the resulting $D^1$-branes has twice the minimal length occurring in the $B_{ab} = 0$ case. The independent branes are thus half of their naive number, and this provides the expected rank reduction, consistently with the doubling of their elementary charge.

In the spirit of section 2, another nice interpretation of the rank reduction for the Chan-Paton group rests on the observation that a quantized $B_{ab}$ is equivalent to a toroidal shift-orbifold compactification. Let us consider for simplicity a two-torus that is the product of two circles of radius $R$, and let us choose $B_{ab}$ in eq. (1) so that the two-dimensional blocks are parametrized as

$$B = \frac{\alpha'}{2} \begin{pmatrix} 0 & 1 \\ -1 & 0 \end{pmatrix}. \quad (3)$$

The non-trivial background makes the sum over momenta dependent on the parity of the integer windings, in such a way that the one-loop torus partition function can be decomposed in the form

$$\Lambda(B) = \Lambda(m_1, m_2, 2n_1, 2n_2) + \Lambda(m_1 + 1/2, m_2, 2n_1, 2n_2 + 1) + \Lambda(m_1, m_2 + 1/2, 2n_1 + 1, 2n_2) + \Lambda(m_1 + 1/2, m_2 + 1/2, 2n_1 + 1, 2n_2 + 1). \quad (4)$$
where $\Lambda(m_1, m_2, n_1, n_2)$ denotes the two-dimensional lattice sum over momenta $m_a/R$ and windings $n^a R$. This expression is indeed the partition function of an asymmetric $p_1 w_2$ shift-orbifold of a two-torus with $B_{ab} = 0$, provided the radius along the first direction is doubled ($R \rightarrow 2R$). As a result, the toroidal compactification on the two-dimensional torus without vector structure is equivalent to a $p_1 w_2$ shift-orbifold compactification, but on a different torus, characterized by a fundamental cell of double length along the first direction. In order to exhibit the expected rank reduction, it is useful to study the $\Omega$-orientifold after performing a T-duality along the vertical direction, that renders the $p_1 w_2$ shift-orbifold action equivalent to a $p_1 p_2$ shift-orientifold where the projection is combined with an inversion. The unoriented truncation forces $m_2$ and $n_1$ to vanish in the Klein-bottle amplitude, while $m_1$ and $n_2$ vanish along the tube, as demanded by the properties of the corresponding boundary states. As a result, the Klein-bottle amplitude takes the form

$$K = \frac{1}{2} P_1 W_2,$$

where $P_i$ and $W_i$ are the usual one-dimensional momentum and winding sums, obtained as unoriented slices of the Narain lattice, while the trans-
verse annulus amplitude
\[ \tilde{A} = \frac{2^{-3}}{2} N^2 \frac{v_1}{v_2} W_1^e P_2^e , \] (6)
contains sums over even momenta and windings, weighted with the normalization needed for a proper particle interpretation of the gauge vectors. The direct-channel annulus amplitude, corresponding to the open-string one-loop channel, can be written
\[ A = \frac{1}{2} N^2 \left( P_1 + P_1^{1/2} \right) \left( W_2 + W_2^{1/2} \right) , \] (7)
where \( P_1 \) and \( W_1 \) are again the usual one-dimensional momentum and winding lattice sums, respectively, while \( P_1^{1/2} \) and \( W_1^{1/2} \) are the corresponding shifted ones. The consistent Möbius amplitude, responsible for the unoriented projection, is then
\[ M = -\frac{1}{2} N \left( \hat{P}_1 \hat{W}_2 + \hat{P}_1^{1/2} \hat{W}_2^{1/2} \right) , \] (8)
as can be checked via a \( P \)-transformation of the transverse amplitude
\[ \tilde{M} = -N \frac{v_1}{v_2} \left( \hat{W}_1^e \hat{P}_2^e + (-1)^{m_1+m_2} \hat{W}_1^e \hat{P}_2^e \right) , \] (9)
demanded by eq. (6) and the transverse Klein-bottle amplitude
\[ \tilde{K} = \frac{2^5}{2} \frac{v_1}{v_2} W_1^e P_2^e . \] (10)
The presence of the shifted sums in eqs. (7) and (8) signals quite neatly that one of the shifts, \( p_2 \) in this case, is orthogonal to the \( D_1 \)-branes, parallel to the \( O_1 \)-planes and thus horizontal. As already observed in Section 2, the orthogonal shift is responsible for the doublet structure of the \( D_1 \)-branes, forcing the split into two stacks of \( N/2 \) brane-image doublets. The geometrical situation is illustrated in fig. 3. The consequent rank reduction of the Chan-Paton gauge group can be easily derived from the analysis of the tadpole cancellation conditions, or simply by inspection of eqs. (6), (9) and (10). For instance, the previous construction, if applied to an eight-dimensional toroidal compactification of the type IIB superstring, would result in type I models with an \( SO(16) \) Chan-Paton group and the expected rank reduction by a factor of two.

Because of the direct connection between quantized \( B_{ab} \) and shift-orbifolds, it can happen that some combination of the two makes \( B_{ab} \) ineffective. In addition, a quantized \( B_{ab} \) generically introduces an open-string
Figure 3. $B$-field vs shift-orientifold: a) space filling $D_2$-branes with $B = 0$. b) Doublets of $D_1$-branes in the shift-orientifold.

analog of discrete torsion, whose values discriminate between Orthogonal and Symplectic gauge groups. However, the two choices can be continuously connected, passing through unitary groups via suitable open-string Wilson lines. In orbifolds, the quantized $B_{ab}$ also affects the $\Omega$ eigenvalues of the twist fields, introducing again exotic $O_-$-planes responsible for the different unoriented truncations, according to the constraints dictated by the underlying Conformal Field Theory. At the same time, additional matter multiplets emerge in the gauge sector of the models in order to balance the loss of R-R charge of the $O_-$-planes.

4 Magnetized Branes and Chiral Models

Chirality in four-dimensional string or M-theory inspired vacua is not an easy property to obtain. The reason is, roughly speaking, that the long wavelength modes come unavoidably from higher-dimensional spinors, that generically give rise to vector-like configurations. Chiral models, however, can be built using projections that retain a net number of chiral spinors. The first example of chiral asymmetry in type I vacua was exhibited in a class of orientifolds of the $Z$-orbifold. In those examples, the $D9$-branes sit at singular points of the internal manifold. However, chirality can also be obtained via other mechanisms. Widely studied in the past few years are backgrounds equipped with uniform magnetic fields along some internal directions of the branes.
or, in a T-dual language, magnetic fields couple solely to the ends of open strings, interpolating between Dirichlet and Neumann boundary conditions. In compact manifolds, the Dirac quantization condition obeyed by the magnetic charges is equivalent to a rational (non-dense) wrapping of the magnetized D-branes around corresponding p-cycles, counted exactly by the degeneracy of the Landau levels. As a result, a fine-tuning of the angles or of the background field-strengths to (anti-)self-dual configurations allows to project away tachyonic modes and to restore supersymmetry, generically broken by the couplings of different spins to the magnetic field. Moreover, due to the non-vanishing topological charge, magnetized branes transmute to uncharged branes of diverse dimensionalities, i.e. they couple to R-R p-forms of different p, via the Wess-Zumino terms of ref.\textsuperscript{29}. Magnetized branes, thus, contribute simultaneously to tadpole cancellation conditions of branes of different dimensionalities. Typically, a “fat” magnetic brane can replace a stack of localized lower-dimensional branes, reducing the rank of the Chan-Paton group and simultaneously yielding matter-field families determined by the Landau-level degeneracies\textsuperscript{115}. In the presence of at least two kinds of (intersecting) D-branes and of compatible shifts, it is possible to obtain four-dimensional chiral models. As an example, we shall discuss the magnetized orientifold of the \textit{w}_1\textit{w}_2\textit{p}_3 model in table\textsuperscript{1}, but an exhaustive description of all the magnetized orientifolds of the models in table\textsuperscript{1} can be found in ref.\textsuperscript{6}. The unoriented closed spectra of the \textit{w}_1\textit{w}_2\textit{p}_3 model contain \(N = 1\) supergravity coupled to \(7\) chiral multiplets. The tadpole cancellation conditions,

\[
\begin{align*}
n + m + \bar{m} &= 8\ 2^{-\frac{r}{2}} \\
d_1 + 2^{-\frac{r_1}{2}} - 2^{-\frac{r_2}{2}} |k_2 k_3| (m + \bar{m}) &= 8\ 2^{-\frac{r}{2}} \\
d_2 &= 8\ 2^{-\frac{r}{2}}
\end{align*}
\]

show how the open sector contains both uncharged and magnetized D9-branes, whose numbers are labelled by \(n\) and \(m\), together with two sets of D5 branes counted by \(d_i\), \(i = 1, 2\). The \(k_i\) in eq.\textsuperscript{11} are the integer Landau level degeneracies of the internal magnetic field that lies along the second and third tori, while \(r = r_1 + r_2 + r_3\) is the total rank of \(B_{ab}\), whose three two-by-two sub-blocks have ranks \(r_j\). The gauge groups can be chosen to be orthogonal or symplectic for the uncharged branes, while on the magnetized D9-branes the gauge group is unitary. The chiral open spectra are reported in table\textsuperscript{2} where \(C\) denotes chiral multiplets, \(\eta\) is a free sign while \(A\), \(F\), \(S\) and \(Adj\) denote the Antisymmetric, Symmetric, Fundamental and Adjoint representation of the gauge group, respectively. Adding open-string Wilson lines
Table 2. Open spectra of the $w_1w_2p_3$-orientifold models.

| Mult. | Number | Reps. $(n, d_1, d_2, m)$ |
|-------|--------|--------------------------|
| $C$   | $3$    | $(\text{Adj}, 1, 1, 1), (1, \text{Adj}, 1, 1)$ | $(1, 1, \text{Adj})$, $(1, 1, 1, \text{Adj})$ |
| $C_L$ | $2^{2^2+r_3} 2 |k_2 k_3|$ | $(F, 1, 1, F + F)$ |
| $C_L$ | $2^{2^2+r_3} 4 |k_2 k_3|$ | $(1, 1, A)$ |
| $C_L$ | $2^{2^2+r_3} 4 |k_2 k_3|$ | $(1, 1, S)$ |
| $C_L$ | $2^{2^2+r_3} 4 |k_2 k_3|$ | $(1, 1, A)$ |
| $C_L$ | $2^{2^2+r_3} 4 |k_2 k_3|$ | $(1, 1, S)$ |

and brane-antibrane pairs, it is possible to obtain models whose spectra are close to (some extensions of) the Standard Model, and examples of this sort are reviewed in other contributions to this volume. This is certainly a step forward in the understanding of the possible phenomenological implications of a brane-world-like scenario, even if the study of the dynamical stability of this kind of vacua and the lack of a “vacuum selection principle” are still serious open problems.

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