Open-string Models with Broken Supersymmetry

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Abstract

I review the salient features of three classes of open-string models with broken supersymmetry. These suffice to exhibit, in relatively simple settings, the two phenomena of "brane supersymmetry" and "brane supersymmetry breaking". In the first class of models, to lowest order supersymmetry is broken both in the closed and in the open sectors. In the second class of models, to lowest order supersymmetry is broken in the closed sector, but is *exact* in the open sector, at least for the low-lying modes, and often for entire towers of string excitations. Finally, in the third class of models, to lowest order supersymmetry is *exact* in the closed (bulk) sector, but is broken in the open sector. Brane supersymmetry breaking provides a natural solution to some old difficulties met in the construction of open-string vacua.

1. Broken supersymmetry and type-0 models

In this talk I would like to review the key features of some open-string models with broken supersymmetry constructed in [1, 2, 3, 4]. These models may be derived in a systematic fashion from corresponding models of oriented closed strings [5], and once more display a surprising richness compared to them. Since the relevant techniques have been discussed at length in the original papers, I will not present any explicit derivations. Rather, referring to some of the resulting vacuum amplitudes, I will try to illustrate how supersymmetry can be broken at tree level in the bulk, on some branes or everywhere.

Closed-string models with broken supersymmetry were among the first new examples considered in the last decade. In particular, the type-0 models [6] provided the first non-trivial instances of modified GSO projections compatible with modular invariance. In order to describe their partition functions, I will begin by introducing some notation that will be used repeatedly in the following, defining the four level-one SO(8) characters

$$O_8 = \frac{\vartheta_3^4 + \vartheta_4^4}{2\eta^4} \quad , \quad V_8 = \frac{\vartheta_3^4 - \vartheta_4^4}{2\eta^4} \quad ,$$
$$S_8 = \frac{\vartheta_2^4 - \vartheta_1^4}{2\eta^4} \quad , \quad C_8 = \frac{\vartheta_2^4 + \vartheta_1^4}{2\eta^4} \quad , \tag{1}$$

where the ϑ_i are Jacobi theta functions and η is the Dedekind function. In terms of these characters, and leaving aside the contribution of the eight transverse bosonic coordinates, the type II models are described by

$$\mathcal{T}_{IIA} = (V_8 - S_8)(\bar{V}_8 - \bar{C}_8) \quad , \tag{2}$$
$$\mathcal{T}_{IIB} = |V_8 - S_8|^2 \quad , \tag{3}$$

while the type-0A and type-0B models are described by

$$\mathcal{T}_{0A} = |O_8|^2 + |V_8|^2 + S_8\bar{C}_8 + C_8\bar{S}_8 \quad , \tag{4}$$

$$\mathcal{T}_{0B} = |O_8|^2 + |V_8|^2 + |S_8|^2 + |C_8|^2 \quad . \tag{5}$$

In these expressions, the characters (O_8, V_8, S_8, C_8) depend on $q = exp(i2\pi\tau)$, with τ the modulus of the torus, while their conjugates depend of $\bar{q}$. All these characters have power series expansions of the type

$$\chi(q) = q^{h-c/24} \sum_{n=0}^{\infty} d_n\, q^n \quad , \tag{6}$$

where the d_n are integers. The low-lying spectra, essentially manifest in this notation, include in all cases the universal triple $(g_{\mu\nu},\ B_{\mu\nu},\ \phi)$. The corresponding states fill a generic transverse matrix, the direct product of the ground states of the V_8 module and of its conjugate $\bar{V}_8$. In addition, the type-IIA superstring has a Majorana gravitino and a Majorana spinor from the NS-R and R-NS sectors and a vector and a three-form from the R-R sector. The fermions result from pairs of Majorana-Weyl spinors of opposite chiralities, ground states of $V_8\bar{C}_8$ and $S_8\bar{V}_8$, while the nature of the R-R bosons is determined by the direct product of the two inequivalent spinor representations of SO(8). These are the ground states of S_8 and $\bar{C}_8$, and the product $8_s \times 8_c$ indeed decomposes into a vector and a three-form. A similar reasoning shows that the type-IIB superstring has a complex Majorana-Weyl gravitino and a complex Weyl fermion from the NS-R and R-NS sectors, and a scalar, a two-form and a self-dual four-form from the R-R sector. The spectra of the type-0 models are purely bosonic, and can be essentially deduced from these. Aside from the universal triple $(g_{\mu\nu},\ B_{\mu\nu},\ \phi)$, their low-lying excitations include a tachyon, the ground state of the $O_8\bar{O}_8$ sector, while their R-R sectors are two copies of the previous ones, and include a pair of vectors and a pair of three-forms for the 0A model, and a pair of scalars, a pair of two-forms and an unconstrained four-form for the 0B model. Both type-0 models are clearly non-chiral, and are thus free of gravitational anomalies.

Let us now turn to the open descendants of the type-0 models. Their structure is essentially determined by the Klein bottle projection. Leaving aside the contributions of the transverse bosons, the conventional choice, originally discussed in [1], corresponds to

$$\mathcal{K}_{0A} \;=\; \frac{1}{2}(O_8 + V_8) \quad , \tag{7}$$

$$\mathcal{K}_{0B} \;=\; \frac{1}{2}(O_8 + V_8 - S_8 - C_8) \quad . \tag{8}$$

It eliminates the NS-NS two-form $B_{\mu\nu}$, but does not affect the tachyon. The effect of $\mathcal{K}$ can be simply summarized recalling that a *positive* (*negative*) sign implies a *symmetrization* (*antisymmetrization*) of the sectors fixed under left-right interchange. This unoriented projection is frequently called Ω. For instance, $\mathcal{K}_{0B}$ symmetrizes the two NS-NS sectors, described by $|O_8|^2$ and $|V_8|^2$, and antisymmetrizes the two R-R sectors, eliminating the NS-NS $B_{\mu\nu}$ and leaving only a pair of R-R two-forms. In addition, the projected spectrum generally includes invariant combinations of all pairs of sectors interchanged by Ω, that do not contribute to $\mathcal{K}$. Thus, the low-lying spectrum of the projected 0A model includes also a R-R vector and a R-R three-form. As usual, the characters in these direct-channel Klein-bottle amplitudes depend on $q\bar{q} = exp(-4\pi\tau_2)$.

The open sector of the 0A model, described by [1]

$$\mathcal{A}_{0A} \;=\; \frac{n_B^2 + n_F^2}{2}(O_8 + V_8) \;-\; n_B n_F (S_8 + C_8) \quad , \tag{9}$$

$$\mathcal{M}_{0A} \;=\; -\,\frac{n_B + n_F}{2}\hat{V}_8 \;-\; \frac{n_B - n_F}{2}\hat{O}_8 \quad , \tag{10}$$

is not chiral and involves *two* different "real" charges, corresponding to orthogonal or symplectic groups. These enter the partition functions via the dimensions, here n_B and n_F, of the corresponding fundamental representations, and in general are subject to linear relations originating from (massless) tadpole conditions. The low-lying modes are simply identified from the contributions of $\mathcal{A}$ and $\mathcal{M}$. For instance, the model contains two sets of $n_B(n_B-1)/2$ and $n_F(n_F-1)/2$ vectors (corresponding to V_8), enough to fill the adjoint representations of a pair of orthogonal groups, tachyons (corresponding to O_8) in doubly (anti)symmetric representations and fermions (corresponding to the R characters S_8 and C_8) in bi-fundamental representations. This description of open-string spectra is also useful in Conformal Field Theory, where it provides a convenient encoding of the spectrum of boundary operators in a generating function of their multiplicities. In this case, if one insists on demanding the cancellation of all NS-NS tadpoles, the result is the family of gauge groups SO(n_B)× SO(n_F), with $n_B+n_F=32$. Here $\mathcal{A}$ depends on $(q\bar{q})^{1/4}=exp(-\pi\tau_2)$. On the other hand, $\mathcal{M}$ depends on $-(q\bar{q})^{1/4}=exp(-\pi\tau_2+i\pi)$, but the "hatted" characters are redefined by suitable phases, and are thus real.

The open sector of the 0B model, described by [1]

$$\begin{aligned}
\mathcal{A}_{0B} &= \frac{n_o^2+n_v^2+n_s^2+n_c^2}{2}V_8 \;+\; (n_o n_v+n_s n_c)O_8 \\
&- (n_v n_s+n_o n_c)S_8 \;-\; (n_v n_c+n_o n_s)C_8 \quad , \qquad (11)\\
\mathcal{M}_{0B} &= -\frac{n_o+n_v+n_s+n_c}{2}\hat{V}_8 \quad , \qquad (12)
\end{aligned}$$

involves four different "real" charges, and is *chiral* but free of anomalies, as a result of the R-R tadpole conditions $n_o=n_v$ and $n_s=n_c$. All irreducible gauge and gravitational anomalies cancel as a result of the R-R tadpole conditions, while the residual anomaly polynomial requires a generalized Green-Schwarz mechanism [26, 2]. If one insists on demanding the cancellation of all NS-NS tadpoles, not related to anomalies [8] as the previous ones, one obtains the family of gauge groups SO(n_o)× SO(n_v)×SO(n_s)× SO(n_c), with $n_o+n_v+n_s+n_c=64$. Despite their apparent complication, these open-string models are actually simpler than the previous ones, since the 0B torus amplitude corresponds to the "charge-conjugation" modular invariant. This circumstance implies a one-to-one correspondence between types of boundaries and types of bulk sectors typical of the "Cardy case" of boundary CFT [9]. In equivalent terms, this model has four types of boundary states, in one-to-one correspondence with the chiral sectors of the bulk spectrum. The boundary states of the 0A model are a bit subtler, since they are proper combinations of these that do not couple to the R-R states, that cannot flow in the transverse channel compatibly with 10D Lorentz invariance. Indeed, the product of the two spinor representations 8_s and 8_c does not contain the identity, and consequently a right-moving S state cannot reflect into a left-moving C state at a Lorentz-invariant boundary.

The modified Klein bottle projection

$$\mathcal{K}'_{0B}=\frac{1}{2}(-O_8+V_8+S_8-C_8) \quad , \qquad (13)$$

first proposed in [2] as an amusing application of the results of [7], removes the tachyon and the NS-NS two-form $B_{\mu\nu}$ from the closed spectrum, and leaves a *chiral* unoriented closed spectrum that comprises the $(g_{\mu\nu},\phi)$ NS-NS pair, together with a two-form, an additional scalar and a self-dual four-form from the R-R sectors. The resulting open spectrum, described by

$$\begin{aligned}
\mathcal{A}'_{0B} &= -\frac{n^2+\bar{n}^2+m^2+\bar{m}^2}{2}C_8+(n\bar{n}+m\bar{m})V_8 \\
&+ (n\bar{m}+m\bar{n})O_8-(mn+\bar{m}\bar{n})S_8 \quad , \qquad (14)\\
\mathcal{M}'_{0B} &= \frac{m+\bar{m}-n-\bar{n}}{2}C_8 \quad , \qquad (15)
\end{aligned}$$

involves the "complex" charges of a pair of unitary groups, subject to the R-R tadpole constraint $m-n=32$, that eliminates all (non-abelian) gauge and gravitational anomalies. The notation resorts to pairs of multiplicities [1], say m and $\bar{m}$, to emphasize the different roles of the fundamental and conjugate fundamental representations of a unitary group U(m). The tadpole conditions identify the numerical values of m and $\bar{m}$, but once again one can read the low-lying spectrum directly and conveniently from the amplitudes written in this form. The choice $n=0$ selects a U(32) gauge group, with a spectrum that is free of tachyons both in the closed and in the open sectors. All irreducible (non-abelian) gauge and gravitational anomalies cancel, while the residual anomaly polynomial requires a generalized Green-Schwarz mechanism [26]. On the other hand, the U(1) factor is anomalous, and is thus lifted by a ten-dimensional generalization of the mechanism of [10], so that the effective gauge group of this model is SU(32). Here one does not have the option of eliminating all the NS-NS tadpoles, and as a result a dilaton potential is generated.

These models have also been studied in some detail in [11], first with the aim of connecting them to the bosonic string, and more recently with the aim of relating them to (non-supersymmetric) reductions of M theory. This last approach goes beyond the perturbative analysis, and therefore has the potential of discriminating between the various options. According to [11], both types of $0B$ descendants admit a non-perturbative definition, while the $0A$ descendants do not. It would be interesting to take a closer look at these relatively simple model and try to elicit some manifestation of this phenomenon.

Let us now spend a few words to summarize the key features of these descendants, where supersymmetry is broken both in the closed and in the open sectors. Whereas the first two models have tachyons both in the closed and in the open sectors, the last results from a non-tachyonic brane configuration of impressive simplicity. This feature actually extends to lower-dimensional compactifications, as first shown by Angelantonj [12]. These type 0 models, and in particular the non-tachyonic one, have interesting applications [13] in the framework of the AdS/CFT correspondence [14].

2. Scherk-Schwarz deformations and brane supersymmetry

We may now turn to the second class of models. These rest on elegant extensions of the Kaluza-Klein reduction, known as Scherk-Schwarz deformations [15], that allow one to induce the breaking of supersymmetry from the different behaviors of fermionic and bosonic modes in the internal space. This setting, as adapted to the entire perturbative spectra of models of oriented closed strings in [16], is the starting point for the constructions in [17, 3]. I will confine my attention to particularly simple examples, related to the reduction of the type IIB superstring on a circle of radius R where the momenta or the windings are subjected to 1/2-shifts, compatibly with modular invariance, in such a way that all massless fermions are lifted in mass. In these models, supersymmetry is completely broken, but several more complicated open-string models, with partial breaking of supersymmetry are discussed in [3, 20].

This Scherk-Schwarz deformation generically introduces tachyons, in the first case *(momentum shifts)* for $R<\sqrt{\alpha'}$, and in the second case *(winding shifts)* for $R>\sqrt{\alpha'}$. The former choice [17, 3] is essentially a Scherk-Schwarz deformation of the low-energy field theory, here lifted to the entire string spectrum. On the other hand, the latter [3] is a bit more subtle to interpret from a field theory perspective, and indeed the resulting deformation of the spectrum is removed in the limit of small radius R, that strictly speaking is inaccessible to the field theory description. Naively, in the first case the open descendants should not present new subtleties. One would expect that the momentum deformations are somehow inherited by the open spectrum, and this is indeed what happens. On the other hand, naively the open spectrum should be insensitive to winding deformations, simply because the available Neumann strings have only momentum excitations. Here the detailed analysis settles the issue in an interesting way. The open spectrum is indeed affected, although in a rather subtle way, and supersymmetry is effectively broken again, at the compactification scale $1/R$, but is *exact* for the massless modes. In order to appreciate this result, let us present the closed-string amplitudes for the two cases,

$$\begin{aligned} \mathcal{T}_1 &= Z_{m,2n}(V_8\bar{V}_8 + S_8\bar{S}_8) + Z_{m,2n+1}(O_8\bar{O}_8 + C_8\bar{C}_8) \\ &- Z_{m+1/2,2n}(V_8\bar{S}_8 + S_8\bar{V}_8) - Z_{m+1/2,2n+1}(O_8\bar{C}_8 + C_8\bar{O}_8) \end{aligned} \tag{16}$$

and

$$\begin{aligned} \mathcal{T}_2 &= Z_{2m,n}(V_8\bar{V}_8 + S_8\bar{S}_8) + Z_{2m+1,n}(O_8\bar{O}_8 + C_8\bar{C}_8) \\ &- Z_{2m,n+1/2}(V_8\bar{S}_8 + S_8\bar{V}_8) - Z_{2m+1,n+1/2}(O_8\bar{C}_8 + C_8\bar{O}_8) \quad , \end{aligned} \tag{17}$$

and the corresponding Klein bottle projections

$$\mathcal{K}_1 = \frac{1}{2}\,(V_8 - S_8)\;Z_m \ , \tag{18}$$

$$\mathcal{K}_2 = \frac{1}{2}\,(V_8 - S_8)\;Z_{2m} + \frac{1}{2}\,(O_8 - C_8)\;Z_{2m+1} \quad . \tag{19}$$

In these expressions, $Z_{m,n}$ denotes the usual Narain lattice sum for the circle

$$Z_{m,n} = \sum_{m,n} q^{\frac{\alpha'}{4}(\frac{m}{R}+\frac{nR}{\alpha'})^2}\;\bar{q}^{\frac{\alpha'}{4}(\frac{m}{R}-\frac{nR}{\alpha'})^2} \quad , \tag{20}$$

while, for instance, $Z_{2m,n}$ denotes the sum restricted to even momenta. In a similar fashion, Z_m in (19) denotes the restriction of the sum to the momentum lattice.

In writing the corresponding open sectors, I will now eliminate several contributions, restricting the charge configurations in such a way that no tachyons are introduced. This is, to some extent, in the spirit of the previous discussion of the 10D U(32) model, but here one can also cancel all NS-NS tadpoles. Moreover, I will take into account the infrared subtlety discussed in [18], that in the model with winding shifts leads the emergence of additional tadpoles in the singular limit $R \to 0$, where whole towers of massive excitations collapse to zero mass. With this proviso, the corresponding open spectra are described by

$$\begin{aligned} \mathcal{A}_1 &= \frac{n_1^2+n_2^2}{2}(V_8 Z_m - S_8 Z_{m+1/2}) + n_1 n_2 (V_8 Z_{m+1/2} - S_8 Z_m) \quad , \\ \mathcal{M}_1 &= -\frac{n_1+n_2}{2}(\hat{V}_8 Z_m - \hat{S}_8 Z_{m+1/2}) \quad , \end{aligned} \tag{21}$$

and

$$\begin{aligned} \mathcal{A}_2 &= \frac{n_1^2+n_2^2}{2}(V_8 - S_8) Z_m + n_1 n_2 (O_8 - C_8) Z_{m+1/2} \quad , \\ \mathcal{M}_2 &= -\frac{n_1+n_2}{2}\left(\hat{V}_8 - (-1)^m \hat{S}_8\right) Z_m \quad , \end{aligned} \tag{22}$$

As anticipated, the first model, with $n_1 + n_2 = 32$, is essentially a conventional Scherk-Schwarz deformation of the type-I superstring. It can also describe the type I spectrum at a finite temperature related to the internal radius R. On the other hand, the second model, where $n_1 = n_2 = 16$, is more interesting, and displays the first novel phenomenon reviewed here, "brane supersymmetry": although supersymmetry is broken at the compactification scale by the Scherk-Schwarz deformation, the massless modes of the open sector fill complete supersymmetry multiplets. I would like to stress that the breaking of supersymmetry in the massive spectrum of the second model can also be regarded as a deformation, now resulting from the unpairing of the Chan-Paton representations for bosonic and fermionic modes with lattice excitations at alternate massive levels. This is the simplest instance of the phenomenon that, following [19], we can call "brane supersymmetry". Here the residual supersymmetry is present only for the massless modes, but in more complicated models it extends to entire sectors of the open spectrum, as first shown in [20].

T-dualities turn these descriptions into equivalent ones, that can often have more intuitive appeal [21]. This is particularly rewarding in the second case: a T-duality along the circle can turn the winding deformation into a momentum deformation orthogonal to the brane responsible for the open-string excitations. It is then perhaps simpler to accept the previous result that the deformation, now a momentum shift orthogonal to the brane, does not affect its massless excitations. A little more work results in a duality argument that associates the (now momentum) shift to the eleventh dimension of M theory, thus realizing the proposal of [22]. Thus, as is often the case, a simple perturbative type I phenomenon has a non-perturbative origin in the heterotic string (and vice versa).

3. Brane supersymmetry breaking

I will now conclude by reviewing the third possibility afforded by these constructions. Here I will follow [3], concocting a six-dimensional analogue of "discrete torsion" [23]. The construction of the resulting closed string model is another application of the methods of [7], in the same spirit as the construction of the 10D U(32) model. Starting from the T^4/Z_2 U(16) $\times$ U(16) model [24, 25], one can revert the Klein-bottle projection for all twisted states. This results in an unoriented closed spectrum with (1,0) supersymmetry, whose massless excitations, aside from the gravitational multiplet, comprise 17 tensor multiplets and 4 hypermultiplets. In [4] it is shown how this choice, described by

$$\begin{aligned}
\mathcal{T} &= \frac{1}{2}|Q_o+Q_v|^2\Lambda + \frac{1}{2}|Q_o-Q_v|^2\left|\frac{2\eta}{\theta_2}\right|^4 \\
&+ \frac{1}{2}|Q_s+Q_c|^2\left|\frac{2\eta}{\theta_4}\right|^4 + \frac{1}{2}|Q_s+Q_c|^2\left|\frac{2\eta}{\theta_3}\right|^4 \quad ,
\end{aligned} \tag{23}$$

$$\mathcal{K} = \frac{1}{4}\left\{(Q_o+Q_v)(P+W) - 2\times 16(Q_s+Q_c)\right\} \quad , \tag{24}$$

does not allow a consistent supersymmetric solution of the tadpole conditions. A consistent solution does exist [4], but requires the introduction of anti-branes, with the end result that supersymmetry, exact to lowest order in the bulk, is necessarily broken on their world volume. Hence the name "brane supersymmetry breaking" for this peculiar phenomenon, that has the attractive feature of confining the breaking of supersymmetry, and the resulting contributions to the vacuum energy, to a brane, or to a collection of branes, that float in a bath of supersymmetric gravity. In writing these expressions, I have introduced the $(1,0)$ supersymmetric characters [1]

$$\begin{aligned}
Q_o &= V_4O_4 - C_4C_4 \quad , \qquad Q_v = O_4V_4 - S_4S_4 \quad , \\
Q_s &= O_4C_4 - S_4O_4 \quad , \qquad Q_c = V_4S_4 - C_4V_4 \quad .
\end{aligned} \tag{25}$$

The two untwisted ones, Q_o and Q_v, start with a vector multiplet and a hypermultiplet, and are Z_2 orbifold breakings of $(V_8 - S_8)$. Out of the two twisted ones Q_s and Q_c, only Q_s describes massless modes, that in this case correspond to a half-hypermultiplet. The breaking of supersymmetry is demanded by the consistency of String Theory. This can be seen rather neatly from the dependence of the transverse-channel Klein bottle amplitude at the origin of the lattices on the sign ϵ associated to the twisted states

$$\tilde{\mathcal{K}}_0 = \frac{2^5}{4}\left\{ Q_o \left(\sqrt{v} \pm \frac{1}{\sqrt{v}}\right)^2 + Q_v \left(\sqrt{v} \mp \frac{1}{\sqrt{v}}\right)^2 \right\} , \tag{26}$$

where the upper signs would correspond to the conventional case of the U(16) $\times$ U(16) model, while the lower signs correspond to the model of eq. (24). Since the terms with different powers of $\sqrt{v}$ are related by tadpole conditions to the multiplicities of the N and D charge spaces, a naive solution of the model corresponding to the upper signs would require a negative multiplicity, $D = -32$.

The open sector is described by

$$\begin{aligned}
\mathcal{A} &= \frac{1}{4}\Bigg\{ (Q_o + Q_v)(N^2 P + D^2 W) + 2ND(Q'_s + Q'_c)\left(\frac{\eta}{\theta_4}\right)^2 \\
&+ (R_N^2 + R_D^2)(Q_o - Q_v)\left(\frac{2\eta}{\theta_2}\right)^2 + 2R_N R_D(-O_4 S_4 - C_4 O_4 + V_4 C_4 + S_4 V_4)\left(\frac{\eta}{\theta_3}\right)^2 \Bigg\} \\
\mathcal{M} &= -\frac{1}{4}\Bigg\{ NP(\hat{O}_4\hat{V}_4 + \hat{V}_4\hat{O}_4 - \hat{S}_4\hat{S}_4 - \hat{C}_4\hat{C}_4) - DW(\hat{O}_4\hat{V}_4 + \hat{V}_4\hat{O}_4 + \hat{S}_4\hat{S}_4 + \hat{C}_4\hat{C}_4) \\
&- N(\hat{O}_4\hat{V}_4 - \hat{V}_4\hat{O}_4 - \hat{S}_4\hat{S}_4 + \hat{C}_4\hat{C}_4)\left(\frac{2\hat{\eta}}{\hat{\theta}_2}\right)^2 + D(\hat{O}_4\hat{V}_4 - \hat{V}_4\hat{O}_4 + \hat{S}_4\hat{S}_4 - \hat{C}_4\hat{C}_4)\left(\frac{2\hat{\eta}}{\hat{\theta}_2}\right)^2 \Bigg\} .
\end{aligned} \tag{27}$$

Supersymmetry is broken on the antibranes, and indeed the amplitudes involve new characters Q'_s and Q'_c, that describe supermultiplets of a chirally flipped supercharge and may be obtained from eq. (25) upon the interchange of S_4 and C_4, as well as other non-supersymmetric combinations. The tadpole conditions determine the gauge group [SO(16) $\times$ SO(16) $]_9 \times$ [USp(16) $\times$ USp(16) $]_{\bar{5}}$, and the 99 spectrum is supersymmetric, with (1,0) vector multiplets for the SO(16) $\times$ SO(16) gauge group and a hypermultiplet in the $(\mathbf{16}, \mathbf{16}, \mathbf{1}, \mathbf{1})$. On the other hand, the $\bar{5}\bar{5}$ spectrum is not supersymmetric and, aside from the [USp(16) $\times$ USp(16)] gauge vectors, contains quartets of scalars in the $(\mathbf{1}, \mathbf{1}, \mathbf{16}, \mathbf{16})$, right-handed Weyl fermions in the $(\mathbf{1}, \mathbf{1}, \mathbf{120}, \mathbf{1})$ and $(\mathbf{1}, \mathbf{1}, \mathbf{1}, \mathbf{120})$ and left-handed Weyl fermions in the $(\mathbf{1}, \mathbf{1}, \mathbf{16}, \mathbf{16})$. Finally, the ND sector, also not supersymmetric, comprises doublets of scalars in the $(\mathbf{16}, \mathbf{1}, \mathbf{1}, \mathbf{16})$ and in the $(\mathbf{1}, \mathbf{16}, \mathbf{16}, \mathbf{1})$, and additional (symplectic) Majorana-Weyl fermions in the $(\mathbf{16}, \mathbf{1}, \mathbf{16}, \mathbf{1})$ and $(\mathbf{1}, \mathbf{16}, \mathbf{1}, \mathbf{16})$. These fields are a peculiar feature of six-dimensional space time, where one can define Majorana-Weyl fermions, if the Majorana condition is supplemented by the conjugation in a pseudoreal representation. All irreducible gauge and gravitational anomalies cancel also in this model, while the residual anomaly polynomial requires a generalized Green-Schwarz mechanism [26].

The radiative corrections in this model are quite interesting, since they convey the (soft) breaking to the gravitational sector. At any rate, the situation in which a model requires the simultaneous presence of branes and presents itself in at least two other instances [27], the four dimensional $Z_2 \times Z_2$ model with discrete torsion and the four-dimensional Z_4 model. In both cases, brane supersymmetry breaking allows a solution of all tadpole conditions and a consistent definition of the open descendants.

One can actually enrich these constructions, allowing for the simultaneous presence of branes and antibranes [29, 28, 4]. These configurations are generically unstable, and their instability reflects itself in the presence of tachyonic excitations, a feature that we have already confronted in our analysis of the ten-dimensional type-0 models. The internal lattice can be used to lift in mass the tachyons, at least within certain ranges of parameters for the internal geometry, that is actually partly *stabilized*, as a result of the different scaling behavior ($O(\sqrt{v})$ and $O(1/\sqrt{v})$) of the contributions of the different D_p branes. Thus, for instance, starting from the type-IIB superstring, one can introduce both branes and anti-branes at the price of having tachyonic excitations in the open spectrum [29]. In addition, even with special tachyon-free configurations, simply waiving the restriction to configurations free of NS-NS tadpoles often gives new interesting models with broken supersymmetry. The simplest setting is provided again by the type-IIB superstring that, aside from the type-I superstring, has an additional chiral tachyon-free descendant, free of gauge and gravitational anomalies, but with broken supersymmetry and a USp(32) gauge group. In lower-dimensional models, more possibilities are afforded by the internal lattice, that may be used to lift in mass some tachyons, leading to stable vacuum configurations including both branes and antibranes. Some of these [30], related to the four-dimensional Z_3 orientifold of [31], appear particularly interesting.

I would like to conclude by mentioning that brane configurations similar to these have also been studied by several other authors over the last couple of years, from a different vantage point, following Sen [32]. Brane studies are reminiscent of monopole studies in gauge theories, and are an interesting enterprise in their own right. Open-string vacua are particular brane configurations that are also vacuum configurations for a perturbative construction, and thus become exact solutions in the limit of vanishing string coupling. This, in retrospect, makes them particularly attractive and instructive, and makes their study particularly rewarding. In addition, they have very amusing applications, in particular to issues related to the AdS/CFT correspondence, that we are only starting to appreciate.

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