CERN LIBRARIES, GENEVA

CM-P00062214

CERN-TH/2003-242

Suppressing the Cosmological Constant in non-Supersymmetric Open-String Vacua

Carlo Angelantonj

Theory Division, CERN — CH-1211 Geneva 23

Abstract

We review the construction of non-supersymmetric type I string models which correspond to consistent flat-space solutions of all classical equations of motion. Moreover, the one-loop vacuum energy is naturally fixed by the size of compact extra dimensions which, in the two-dimensional case, can be lowered to a fraction of a millimetre. This class of models has interesting non-abelian gauge groups and can accommodate chiral fermions. In the large radius limit, supersymmetry is recovered in the bulk, while D-brane excitations, although non-supersymmetric, exhibit Fermi-Bose degeneracy at all mass levels.

SUPPRESSING THE COSMOLOGICAL CONSTANT IN NON-SUPERSYMMETRIC OPEN-STRING VACUA

CARLO ANGELANTONJ

Theory Division, CERN — CH-1211 Geneva 23

E-mail: carlo.angelantonj@cern.ch

We review the construction of non-supersymmetric type I string models which correspond to consistent flat-space solutions of all classical equations of motion. Moreover, the one-loop vacuum energy is naturally fixed by the size of compact extra dimensions which, in the two-dimensional case, can be lowered to a fraction of a millimetre. This class of models has interesting non-abelian gauge groups and can accommodate chiral fermions. In the large radius limit, supersymmetry is recovered in the bulk, while D-brane excitations, although non-supersymmetric, exhibit Fermi-Bose degeneracy at all mass levels.

1 Introduction

One of the outstanding problems in string theory, as in any quantum theory including gravity, is to understand the smallness of the cosmological constant. This problem became much more severe after the recent observations suggesting a non-vanishing value corresponding to a new energy scale, by far smaller than every other scale in physics of fundamental interactions, $\Lambda \equiv E_A^4$ with $E_A \sim 10^{-4}$ eV. On the other hand, brane-world models with large extra dimensions of submillimetre size have recently received a lot of attention, as they offer a new framework for stabilising the mass hierarchy and addressing several problems of particle physics in an interesting new perspective ^{1,2}.

It is then suggestive, if not a simple numerical coincidence, that the size of the bulk $R_\perp$ in these models is of the same order of magnitude as the cosmological constant scale $R_\perp \sim E_A^{-1} \sim 1 \text{ mm}^a$. In fact, the behaviour $\Lambda \sim 1/R_\perp^4$ is generically valid in models where supersymmetry is broken by Scherk-Schwarz boundary conditions ^{4,5,6,7,1,8}. However, in this case the induced mass splittings are of order $1/R_\perp$, and thus too small by several orders of magnitude. It would then be very appealing if supersymmetry might be broken on the branes at the string scale $M_s \sim 1 \text{ TeV}$, in a way that the vacuum energy vanishes in the decompactification limit.

Type II string models with possible vanishing cosmological constant in perturbation theory were studied in refs. ^{9,10}. Their main feature is a Fermi-

^aSee ³ where similar ideas for suppressing the cosmological constant have been exposed, though employing different methods.

Figure 1. O -planes and D-branes configurations for the seven-dimensional orientifold.

Bose degenerate spectrum, thus leading to an automatic vanishing of the one-loop vacuum energy. Aside from the question of possible higher-loop corrections¹¹, their main defect, however, was that the non-abelian gauge sector, appearing at particular singular points of the compactification manifold, or on appropriate D-brane collections, was always supersymmetric^{12,13}. Thus, it is questionable whether such constructions can accommodate gauge degrees of freedom with large supersymmetry-breaking mass splittings.

Here, we review the construction in¹⁴ of type I string models with supersymmetry broken on the D-branes at the string scale M_s , and an additional source of supersymmetry breaking in the bulk induced by Scherk-Schwarz boundary conditions. In the large radius limit, the one-loop cosmological constant vanishes as $1/R_1^4$, while the D-brane spectra exhibit fermion-boson degeneracy at all massive levels. In the simplest case, this degeneracy can be thought of as emerging from a supersymmetric theory, where superpartners have been displaced appropriately in the position space.

2 The Model

Following¹⁴, we present here the simplest instance of a non-supersymmetric orientifold model^b with suppressed one-loop cosmological constant. To this end, one has simply to compactify the nine-dimensional model with M-theory breaking of⁶ on an additional T^2 with a non-vanishing B -field¹⁶, and distribute the D-branes as in figure 1.

To be more concrete, the parent type IIB model is encoded in the torus

^bSee¹⁵ for reviews on (non-)supersymmetric orientifold constructions

amplitude

$$\mathcal{T} = \left[(V_8 \bar{V}_8 + S_8 \bar{S}_8) \Gamma_{m,2n}^{(1,1)} + (O_8 \bar{O}_8 + C_8 \bar{C}_8) \Gamma_{m,2n+1}^{(1,1)} - (V_8 \bar{S}_8 + S_8 \bar{V}_8) \Gamma_{m+\frac{1}{2},2n}^{(1,1)} - (O_8 \bar{C}_8 + C_8 \bar{O}_8) \Gamma_{m+\frac{1}{2},2n+1}^{(1,1)} \right] \Gamma^{(2,2)}(B).$$

To fix the notation, x^9 is the coordinate on the circle of radius R responsible for the Scherk-Schwarz deformation, with m and n its associated momenta and windings. Coordinates on the spectator T^2 are labelled by the pair (x^7, x^8) . For simplicity we choose T^2 to be a squared torus with sides of size $2\pi R_7$ and $2\pi R_8$, and, of course, with a background $B_{ab} = \frac{\alpha'}{2} \epsilon_{ab}$. The corresponding momenta and windings are labelled by (m_7, m_8) and (n^7, n^8) .

Thus, following¹⁷, we can immediately write down the Klein-bottle amplitude

$$\mathcal{K} = \frac{1}{2} \left[(V_8 - S_8) W_{2n}^{(1)} + (O_8 - C_8) W_{2n+1}^{(1)} \right] W_{(2n^7, 2n^8)}^{(2)}$$

associated to the $\Omega\mathcal{J}$ orientifold projection, with now $\mathcal{J}x^i = -x^i$, an inversion on the three compact coordinates, $i = 7, 8, 9$.

In terms of ten-dimensional fields, the massless closed unoriented spectrum comprises the graviton, the dilaton and a R-R two-form potential. For $R < \sqrt{\alpha'}$ a tachyon appears in the spectrum in the odd-winding (twisted) sector. However, here we are only interested in very large values for R and thus we are not concerned with this.

An S modular transformation maps $\mathcal{K}$ into the transverse-channel amplitude

$$\begin{aligned} \tilde{\mathcal{K}} = & \frac{2^5}{2 \cdot 4} \frac{(\alpha')^{3/2}}{RR_8 R_7} \left[V_8 \left(\frac{1 + (-1)^m}{2} \right)^2 - S_8 \left(\frac{1 - (-1)^m}{2} \right)^2 \right] P_m^{(1)} \\ & \times \left(\frac{1 - (-1)^{m_7} + (-1)^{m_8} + (-1)^{m_7+m_8}}{2} \right)^2 P_{(m_7, m_8)}^{(2)}. \end{aligned}$$

As expected, one recognises the distribution of orientifold planes depicted in figure 1.

The rest of the construction is now straightforward. We have simply to add eight brane-antibrane pairs to cancel the NS-NS tadpole in $\tilde{\mathcal{K}}$. As usual, in the open sector we have the freedom to distribute the branes along the compact dimensions. In the case at hand, the configuration of figure 1 corresponds to the annulus amplitude

$$\mathcal{A} = \frac{1}{2} (N_D^2 + N_{\bar{D}}^2) (V_8 - S_8) W_n^{(1)} W_{(n^7, n^8)}^{(2)} + N_D N_{\bar{D}} (O_8 - C_8) W_{n+\frac{1}{2}}^{(1)} W_{(n^7+\frac{1}{2}, n^8)}^{(2)},$$

where, as before, N_D ($N_{\bar{D}}$) counts the number of branes (antibranes).

From

$$\tilde{\mathcal{A}} = \frac{2^{-5} (\alpha')^{3/2}}{2 R R_8 R_7} \left[((-1)^m (-1)^{m_7} N_D + N_{\bar{D}})^2 V_8 - ((-1)^m (-1)^{m_7} N_D - N_{\bar{D}})^2 S_8 \right] P_m^{(1)} P_{(m_7, m_8)}^{(2)}$$

and $\tilde{\mathcal{K}}$ we can immediately extract the Möbius amplitude in the transverse channel

$$\tilde{\mathcal{M}} = -\frac{1}{2} \frac{(\alpha')^{3/2}}{R R_8 R_7} \left[V_8 \left(\frac{1 + (-1)^m}{2} \right) + S_8 \left(\frac{1 - (-1)^m}{2} \right) \right] P_m^{(1)} \times \left(\frac{1 - (-1)^{m_7} + (-1)^{m_8} + (-1)^{m_7 + m_8}}{2} \right) ((-1)^{m_7} N_D + N_{\bar{D}}) P_{(m_7, m_8)}^{(2)}.$$

Upon a P -modular transformation, one then finds

$$\mathcal{M} = -\frac{1}{2} (V_8 + (-1)^n S_8) W_n^{(1)} \left[(N_D + N_{\bar{D}}) W_{(n^7, 2n^8 + 1)}^{(2)} - (N_D - N_{\bar{D}}) (-1)^{n^7} W_{(n^7, 2n^8)}^{(2)} \right],$$

which is the proper (anti-)symmetrisation of $\mathcal{A}$.

The D-brane massless spectrum thus comprises the dimensional reduction of ten-dimensional gauge bosons with gauge group $\text{USp}(8) \otimes \text{SO}(8)$ and fermions in the $(28, 1) \oplus (1, 36)$ representations. The two gauge group factors (*i.e.* the branes and the antibranes) are located at different points along the compact (x^8, x^9) directions, as depicted in figure 1. This massless spectrum is obviously Fermi-Bose degenerate, although not supersymmetric. Moreover, in the limit $R, R_8 \rightarrow \infty$ all massive string levels present Fermi-Bose degeneracy.

3 The vacuum energy density

The one-loop vacuum energy receives contributions from the four surfaces of vanishing Euler character

$$\Lambda(R) = \frac{1}{2} \int_{\mathcal{F}} \frac{d^2 \tau}{\tau_2^{9/2}} \frac{\mathcal{T}(R)}{|\eta|^{10}} + \int_0^\infty \frac{d\tau_2}{\tau_2^{9/2}} \frac{\mathcal{K}(R)}{\eta^5} + \int_0^\infty \frac{d\tau_2}{\tau_2^{9/2}} \frac{\mathcal{A}(R)}{\eta^5} + \int_0^\infty \frac{d\tau_2}{\tau_2^{9/2}} \frac{\mathcal{M}(R)}{\hat{\eta}^5},$$

where $\mathcal{T}(R)$, $\mathcal{K}(R)$, $\mathcal{A}(R)$ and $\mathcal{M}(R)$ are given in the previous section.

Figure 2. Mass and charge distribution in the $R, R_8 \gg R_7$ limit.

For $\tau_2 > 0$, the contributions from the Klein-bottle, annulus and Möbius-strip amplitudes can be made arbitrarily small by taking two of the compact directions, R and R_8 , very large, since

$$\Lambda(\mathcal{K}) + \Lambda(\mathcal{A}) + \Lambda(\mathcal{M}) \sim \mathcal{O}(e^{-R^2} \text{ or } e^{-R_8^2}).$$

On the other hand, the potentially divergent region $\tau_2 = 0$ is better described, as usual, in terms of the horizontal time $\ell \sim 1/\tau_2$. One can then show that in the $\ell \rightarrow \infty$ limit, and for $R, R_8 \gg R_7$, the vacuum energy density reads

$$\Lambda \simeq \frac{128}{\pi} \frac{R\sqrt{\alpha'}}{R_7 R_8} \sum \frac{(-1)^m - (-1)^{m_8}}{m^2 + \left(\frac{R}{R_8}\right)^2 m_8^2},$$

and vanishes in the symmetric case $R = R_8$, up to exponentially small corrections, as stated before.

Actually, these amplitudes can be given a suggestive geometrical interpretation in terms of gravitational and “Maxwell”-type interactions¹⁸, where the condition $R = R_8$ would translate into an overall compensation of forces (a BPS-like condition). In the limit $R_7 \ll R, R_8$, we are mainly left with a two-dimensional distribution of charged and massive point-like sources as in figure 2. Here, the two entries (M_i, q_i) in the parenthesis refer to the mass and charge of the i -th “particle”, where in our conventions a single $\mathcal{O}7^+$ plane has both tension and charge equal to -4 while a single D-brane carries a (positive) unit of tension and charge. As a result, the overall gravitational force

experienced by the system is

$$\begin{aligned} F_{\text{Newton}} &= -(M_1 M_2 + M_3 M_4) \log R - (M_1 M_4 + M_2 M_3) \log R_8 \\ &\quad - (M_1 M_3 + M_2 M_4) \log \sqrt{R^2 + R_8^2} \\ &= -128 \log R + 128 \log R_8 + 128 \log \sqrt{R^2 + R_8^2}. \end{aligned}$$

Similarly, the “Maxwell” force due to the R-R field exchange reads

$$\begin{aligned} F_{\text{Maxwell}} &= +(q_1 q_2 + q_3 q_4) \log R + (q_1 q_4 + q_2 q_3) \log R_8 \\ &\quad + (q_1 q_3 + q_2 q_4) \log \sqrt{R^2 + R_8^2} \\ &= -128 \log R + 128 \log R_8 - 128 \log \sqrt{R^2 + R_8^2}, \end{aligned}$$

and thus the total force is

$$F_{\text{Newton}} + F_{\text{Maxwell}} = -256 \log R + 256 \log R_8.$$

It is then clear that, as stated previously, the condition $R = R_8$, required to cancel the leading contribution to the cosmological constant, translates into a (BPS-like) zero-force condition on the distribution of D-branes and $\mathcal{O}$ -planes.

The analysis of the torus amplitude is standard. The twisted (odd-winding) sector does not have massless Kaluza-Klein states and thus gives an exponentially suppressed contribution. On the contrary, the untwisted sector contributes with the familiar power-law fall-off $1/R^7$. Upon compactification to four dimensions (*e.g.* on spectator tori) one recovers the expected $1/R^4$ behaviour.

4 Conclusions

We have here reviewed the construction of the first instance of a vacuum configuration with non-supersymmetric non-abelian gauge sector and a naturally small cosmological constant¹⁴. In the simplest case, the model contains two “mirror worlds”, while the degeneracy is due to an exchange of ordinary superpartners on the two branes.

The main question that remains to be addressed concerns, of course, higher-order corrections, that we leave for further investigations.

Acknowledgements

I would like to thank Ignatios Antoniadis for enjoyable collaborations, and the organisers of the SP2004 conference for their kind invitation.

References

1. I. Antoniadis, "A Possible New Dimension at a Few TeV," Phys. Lett. B246 (1990) 377.
2. N. Arkani-Hamed, S. Dimopoulos and G. Dvali, "The Hierarchy Problem and New Dimensions at a Millimetre," Phys. Lett. B429 (1998) 263, [arXiv:hep-ph/9803315] and "Phenomenology, astrophysics and cosmology of theories with sub-millimetre dimensions and TeV scale quantum gravity," Phys. Rev. D59 (1999) 086004 [arXiv:hep-ph/9807344]; I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos and G. R. Dvali, "New dimensions at a millimetre to a Fermi and superstrings at a TeV," Phys. Lett. B436 (1998) 257 [arXiv:hep-ph/9804398].
3. C. P. Burgess, R. C. Myers and F. Quevedo, "A naturally small cosmological constant on the brane?," Phys. Lett. B495 (2000) 384 [arXiv:hep-th/9911164]; Y. Aghababaie, C. P. Burgess, S. L. Parameswaran and F. Quevedo, "Toward a naturally small cosmological constant from branes in 6D supergravity," [arXiv:hep-th/0304256].
4. J. Scherk and J. H. Schwarz, "How To Get Masses From Extra Dimensions," Nucl. Phys. B153 (1979) 61 and "Spontaneous Breaking Of Supersymmetry Through Dimensional Reduction," Phys. Lett. B82 (1979) 60.
5. R. Rohm, "Spontaneous Supersymmetry Breaking In Supersymmetric String Theories," Nucl. Phys. B237 (1984) 553; C. Kounnas and M. Porrati, "Spontaneous Supersymmetry Breaking In String Theory," Nucl. Phys. B310 (1988) 355; S. Ferrara, C. Kounnas, M. Porrati and F. Zwirner, "Superstrings With Spontaneously Broken Supersymmetry And Their Effective Theories," Nucl. Phys. B318 (1989) 75.
6. I. Antoniadis, E. Dudas and A. Sagnotti, "Supersymmetry breaking, open strings and M-theory," Nucl. Phys. B544 (1999) 469 [arXiv:hep-th/9807011]; I. Antoniadis, G. D'Appollonio, E. Dudas and A. Sagnotti, "Partial breaking of supersymmetry, open strings and M-theory," Nucl. Phys. B553 (1999) 133 [arXiv:hep-th/9812118].
7. H. Itoyama and T. R. Taylor, "Supersymmetry Restoration In The Compactified $O(16) \times O(16)$ -Prime Heterotic String Theory," Phys. Lett. B186 (1987) 129.
8. I. Antoniadis, S. Dimopoulos and G. R. Dvali, "Millimetre range forces in superstring theories with weak-scale compactification," Nucl. Phys. B516 (1998) 70 [arXiv:hep-ph/9710204].
9. S. Kachru, J. Kumar and E. Silverstein, "Vacuum energy cancellation in a non-supersymmetric string," Phys. Rev. D59 (1999) 106004 [arXiv:hep-

- th/9807076].
10. J. A. Harvey, "String duality and non-supersymmetric strings," Phys. Rev. D59 (1999) 026002 [arXiv:hep-th/9807213].
 11. S. Kachru and E. Silverstein, "On vanishing two loop cosmological constants in non-supersymmetric strings," JHEP 9901 (1999) 004 [arXiv:hep-th/9810129]; R. Iengo and C. J. Zhu, "Evidence for non-vanishing cosmological constant in non-SUSY superstring models," JHEP 0004 (2000) 028 [arXiv:hep-th/9912074].
 12. R. Blumenhagen and L. Görlich, "Orientifolds of non-supersymmetric, asymmetric orbifolds," Nucl. Phys. B551 (1999) 601 [arXiv:hep-th/9812158].
 13. C. Angelantonj, I. Antoniadis and K. Forger, "Non-supersymmetric type I strings with zero vacuum energy," Nucl. Phys. B555 (1999) 116 [arXiv:hep-th/9904092].
 14. C. Angelantonj and I. Antoniadis, "Suppressing the cosmological constant in non-supersymmetric type I strings," arXiv:hep-th/0307254.
 15. C. Angelantonj and A. Sagnotti, "Open strings," Phys. Rept. 371 (2002) 1 [Erratum-ibid. 376 (2003) 339] [arXiv:hep-th/0204089]; E. Dudas, "Theory and phenomenology of type I strings and M-theory," Class. Quant. Grav. 17 (2000) R41 [arXiv:hep-ph/0006190].
 16. M. Bianchi, G. Pradisi and A. Sagnotti, "Toroidal compactification and symmetry breaking in open string theories," Nucl. Phys. B376 (1992) 365.
 17. A. Sagnotti, "Open Strings And Their Symmetry Groups," Cargese Summer Institute on Non-Perturbative Methods in Field Theory, Cargese, France, 1987 [arXiv:hep-th/0208020]; M. Bianchi and A. Sagnotti, "On The Systematics Of Open String Theories," Phys. Lett. B247 (1990) 517 and "Twist Symmetry And Open String Wilson Lines," Nucl. Phys. B361 (1991) 519; P. Horava, "Strings On World Sheet Orbifolds," Nucl. Phys. B327 (1989) 461.
 18. J. Polchinski, "Dirichlet-Branes and Ramond-Ramond Charges," Phys. Rev. Lett. 75 (1995) 4724 [arXiv:hep-th/9510017].