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D-terms and D-strings in open string models

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Abstract: We study the Fayet-Iliopoulos (FI) D-terms on D-branes in type-II Calabi-Yau backgrounds. We provide a simple worldsheet proof of the fact that, at tree level, these terms only couple to scalars in closed string hypermultiplets. At the one-loop level, the D-terms get corrections only if the gauge group has an anomalous spectrum, with the anomaly cancelled by a Green-Schwarz mechanism. We study the local type-IIA model of D6-branes at SU(3) angles and show that, as in field theory, the one-loop correction suffers from a quadratic divergence in the open string channel. By studying the closed string channel, we show that this divergence is related to a closed string tadpole, and is cancelled when the tadpole is cancelled. Next, we study the cosmic strings that arise in the supersymmetric phases of these systems in light of recent work of Dvali et al. In the type-IIA intersecting D6-brane examples, we identify the D-term strings as D4-branes ending on the D6-branes. Finally, we use $\mathcal{N} = 1$ dualities to relate these results to previous work on the FI D-term of heterotic strings.

Keywords: Superstrings and Heterotic Strings, String Duality, Brane Dynamics in Gauge Theories.

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1. Introduction

In recent years, the study of D-branes and orientifold planes in nontrivial string backgrounds has been brought under a degree of computational control. Within this framework, one may construct a (bewildering) variety of supersymmetric models, some of them with spectra close to the supersymmetric standard model. Perhaps more interestingly, it is via D-brane models combined with magnetic fluxes that the greatest advances have been made towards constructing vacuum states free of light scalar fields [1]–[6].

In order to address both the practical model-building questions as well as the deeper physical issues raised by this recent work, it is important to have a better quantitative and qualitative understanding of the 4d low-energy effective action (and its validity) for such models. Most of the empirically-based puzzles of particle physics — e.g. the hierarchy problem and the cosmological constant problem — are discussed in this arena.

In open string models, the F-terms can be brought under control due to spacetime nonrenormalization theorems, coupled with the fact that such terms can be computed using topological string theory. The D-term couplings are notoriously more difficult to compute. However, in these models the Fayet-Iliopoulos (FI) D-terms can be understood at tree level e.g. [7, 8], and they are particularly important in understanding the vacuum structure of the theory. Furthermore, the D-term part of the potential is central to recent attempts to construct inflationary models in string theory [9]–[16]. If the inflationary potential is dominated by D-terms, one may avoid some of the naturalness issues that plague F-term-driven inflation [17, 18].

There are several related puzzles regarding these terms in type-II D-brane models. First, a field-theory nonrenormalization theorem [19] states that such D-terms get quantum corrections only at one loop order in perturbation theory, and then only if the chiral multiplets have charges Q_i under this U(1) such that $\sum_i Q_i \neq 0$, i.e. if the spectrum is anomalous. In this case, it is known that the FI D-term gets a quadratically divergent one-loop correction proportional to $(\sum_i Q_i)\Lambda_{UV}^2$, where Λ_{UV} is the ultraviolet cutoff. There are string theory constructions with such a spectrum [20, 21], wherein the anomaly is cancelled by the Green-Schwarz mechanism [22, 23]. Our first question is, then, what is the stringy version of the one-loop calculation above? In the heterotic string, when the anomaly is cancelled by a coupling $F_{\mu\nu}B^{\mu\nu}$ of the U(1) gauge field strength to the NS 2-form potential B, the UV divergence is cut off at the string scale by the restriction of the modular integral to a fundamental domain of the modular group [20, 24]. In open string theories, as pointed out in [25], there is no such mechanism for cutting off the integral over open string moduli. The corresponding cylinder amplitude is thus divergent. In the type-I models discussed in [25, 26, 27], the divergence is cancelled by a diagram with a crosscap, which is intrinsically stringy. One can ask how general this story is.[1]

There exist string models [21] for which there is a single charged chiral multiplet [28], with an FI D-term which one can tune by hand, passing between phases of broken and unbroken supersymmetry [29]. This is a generic local model [30] for such transitions in type-II theories (general type-IIB examples will be mirror to this). During such transitions, the underlying closed string background is nonsingular. In such models there is a BPS cosmic string whose tension is proportional to the FI D-term, as studied recently by Dvali and collaborators [31, 32]. An interesting question then arises as to the identity of this cosmic string in string theory. As predicted by the effective field theory, there should be a tensionless string at the phase boundary between broken and unbroken supersymmetry. Yet this boundary has real codimension one and so should not correspond to some brane wrapping a vanishing cycle in the underlying manifold: rather, we expect the D-branes themselves to become singular at this transition [29, 7]. Therefore, we should seek a D-brane configuration which is a string in spacetime, whose tension vanishes as the collection of space-filling D-branes becomes singular.

[1]This question was asked by E. Poppitz and S. Kachru in 1999.

In this paper we will answer these questions. The outline is as follows. In section 2, we show via a worldsheet argument (similar to that in [33]) that at tree level, the FI D-terms do not depend on the closed-string modes which descend from $N = 2$ vector multiplets. In section 3 we study the one-loop contribution to the D-terms for the example of two D6-branes intersecting in $3+1$ dimensions [21], and argue that the divergent contribution cancels when all disc contributions to massless closed string tadpoles vanishes. In section 4, we study the type-IIA realization of the Fayet model with intersecting D6-branes [29, 30], and study the cosmic strings which arise in the supersymmetric phase, corroborating recent work of Dvali et al. and Binetruy et al. [31, 32] in this more nontrivial class of examples. Our identification is similar to a desciption given in [34, 35, 36].[2] In section 5 we conclude by discussing the relation of our results to dual heterotic and M-theory models. An appendix gives the details of the one-loop open-string calculations.

In this work, we are interested in theories for which nontrivial gauge dynamics arises from D-branes. As in [41, 42], we will focus on "local models", by which we mean that we focus on low-energy D-brane dynamics near some region of interest of the geometry. The low-energy dynamics can be captured by placing D-branes in a noncompact background with the same local structure. There are two advantages to this approach. First, the noncompact models are easier to describe and decouple gravity from the problem. Secondly, they may be glued in to a wide variety of compact models. Other sectors with light fields will often be physically separated in the extra dimensions from the sector of interest. From the standpoint of the local model, this will fix the behavior at infinity.

2. Closed string decoupling theorems for the D-brane effective action

For type-II models at tree level, the fact that the closed strings lie in $\mathcal{N} = 2$ supermultiplets controls how they couple to the FI D-terms: in particular it can be shown that specific closed string modes decouple from specific terms in the $4d$ effective action for open strings. In the case of supersymmetric D-branes in type-II string models, it was argued in [33] by surveying known examples [43, 44] that FI D-terms for open string degrees of freedom are *independent* of closed string modes which descend from $\mathcal{N} = 2$ vector multiplets, while superpotential terms for open string chiral scalar multiplets are independent of closed string modes which descend from $\mathcal{N} = 2$ hypermultiplets. These facts are crucial in the categorical description of supersymmetric D-brane configurations [7, 45].

The statement regarding the superpotential was proven via worldsheet techniques in [33]. We can provide an equally simple proof of the decoupling statement for FI D-terms at tree level. The central point is that $\mathcal{N} = 1$ spacetime supersymmetry requires $\mathcal{N} = 2$ worldsheet superconformal symmetry in the open string sector. The $\mathcal{N} = 2$ algebra contains a $\mathrm{U}(1)_R$ affine Lie subalgebra, and the decoupling theorem for FI D-terms is a consequence of $\mathrm{U}(1)_R$ selection rules. It is worth noting that in this context, the existence of the U(1) R-symmetry is weaker than the requirement of unbroken 4d $\mathcal{N} = 1$ supersymmetry.

[2] Recent papers which also study D-term strings in string theory include [37]–[40].

We are interested in the behavior of the FI D-term under a small deformation $\delta\phi$ of the closed string background. Let this deformation be described by the $(-1,-1)$ picture vertex operator $V_{\delta\phi} = e^{-\phi-\tilde{\phi}}\mathcal{O}$, where $\phi, \tilde{\phi}$ are the left- and right-moving bosonized superconformal ghosts (q.v. [46]) and $\mathcal{O}$ is a dimension $(1/2, 1/2)$ operator. We would also like a vertex operator describing the auxiliary field D for a U(1) gauge group. Let the vertex operator for an open string gauge field have a Chan-Paton matrix $s_{aa'}$, where a, a' label the constituent D-branes. One may easily adapt the treatment of the heterotic string in [24, 47] to the open string case, to show that the (0)-picture boundary operator

$$V_{D,aa'} = J_{\mathrm{U}(1)} s_{aa'} \tag{2.1}$$

is the operator for the FI D-term for $\mathrm{U}(1)_a$, and $s_{aa'} = \delta_{aa'}$. In the case of the heterotic string, s is replaced by the left-moving current for the corresponding element of the world-sheet affine Lie algebra associated to the spacetime gauge group [47]. The variation of the spacetime coupling $\int d^4x \xi D$ with respect to the deformation $\delta\phi$ is simply the disc amplitude:

$$\langle V_D(w_0) V_{\delta\phi}(z, \bar{z}) \rangle \tag{2.2}$$

where $(z, \bar{z})$ is a fixed point in the interior of the disc D and w_0 is a location on the boundary ∂D of the disc.[3]

In the models of interest, the closed string sector has $\mathcal{N} = (2,2)$ worldsheet supersymmetry, and the spacetime fields corresponding to massless vertex operators lie in 4d $\mathcal{N} = 2$ spacetime supermultiplets. If $\delta\phi$ denotes a scalar in a vector multiplet, the corresponding NS-NS vertex operator $\mathcal{O}$ lies in one of the four chiral rings. In type-IIB string theory, these are complex structure deformations with $\mathrm{U}(1)_R$ charge $(1,1)$ or $(-1,-1)$, while in type-IIA string theory they are Kähler deformations with $\mathrm{U}(1)_R$ charge $(\pm 1, \mp 1)$. We are interested in D-branes which fill the 4d spacetime and preserve supersymmetry. As described in [33], the corresponding open string boundary conditions preserve the diagonal combination $J + \tilde{J}$ of left- and right-moving $\mathrm{U}(1)_R$ current algebra for type-IIB, and the off-diagonal combination $J - \tilde{J}$ for type-IIA backgrounds. Therefore $V_{\delta\phi}$ is charged under the preserved $\mathrm{U}(1)_R$ current, while V_D is that R-current itself, and therefore neutral. Therefore eq. (2.2) vanishes by $\mathrm{U}(1)_R$-charge conservation.

Although it is a digression from the theme of this work, we note that the gauge coupling for open string fields also decouples from closed string vectormultiplets at tree level, as the proof is identical to that for the FI D-terms. The variation of the gauge coupling with respect to $\delta\phi$ is proportional to the disc amplitude

$$\langle V_{\delta\phi}(z, \bar{z}) V_A(w_0) \oint_{\partial D} dw V_A(w) \rangle \tag{2.3}$$

where V_A is the $(0,0)$-picture vertex operator for the gauge field, and w_0, w lie on the boundary of the disc. Again, V_A is neutral with respect to the preserved $\mathrm{U}(1)_R$ symmetry of the theory, so this amplitude vanishes by worldsheet R-charge conservation.

[3] V_D is not a vertex operator for a physical state, and the reader may feel more comfortable measuring instead the coupling of $\delta\phi$ to the masses of scalar fields charged under the corresponding U(1), as in [43, 25]. However, an analogous analysis to that of [24, 47] shows that the calculation described here is equivalent.

3. One-loop FI contributions for intersecting D-branes

Consider a general $d = 4, N = 1$ gauge theory with group $G = \mathrm{U}(1) \times G'$. If the chiral multiplets Φ^i have charges q^i, then there is a quadratically-divergent one-loop contribution to the the FI term of the form [19]:

$$\xi_{\mathrm{one-loop}} \propto \left(\sum_i q^i \right) \Lambda_{UV}^2 \tag{3.1}$$

where Λ_{UV} is an ultraviolet cutoff. In other words, the FI term has a divergent contribution when the U(1) gauge symmetry is anomalous.

The standard lore is that in consistent string backgrounds, divergences are either cut off or have an infrared interpretation, and gauge anomalies are cancelled via the Green-Schwarz mechanism. The status of the one-loop contribution to the FI D-terms seems to depend on the model at hand. In the case of the heterotic string, when the anomaly is cancelled by a coupling to the "universal" axion dual to the NS-NS 2-form, the one-loop contribution is finite [24, 20], with $\Lambda_{UV}^2 = M_s^2$: modular invariance removes the UV-divergent region of the string diagrams.

In the case of open string U(1) gauge groups, there is no modular group to cut off the UV region of the open string loop amplitudes. The divergence must be cancelled or explained. For a particular type-I orentifold vacuum [25], the one-loop correction was shown to vanish identically due to a cancellation between the cylinder and Möbius strip diagrams. It is not clear how general this story might be.

In order to answer this more directly, we examine the one-loop correction to the FI D-term in a noncompact type-IIA example of two D6-branes A, B in $\mathbb{R}^{10}$ intersecting along $\mathbb{R}^4$, with angles chosen such that $\mathcal{N} = 1, d = 4$ SUSY is preserved at tree level [21]. The strings localized at the intersection of these branes lead to an anomalous spectrum for the off-diagonal $\mathrm{U}(1)_-$ generated by the CP matrix $\left(\begin{smallmatrix} 1 & 0 \\ 0 & -1 \end{smallmatrix}\right)$. This is a local model of the realization in [29] of the Fayet model and provides a fairly generic picture of how FI D-terms for open strings arise in type-IIA compactifications with branes. In this section we compute the one-loop contribution to the FI D-term for $\mathrm{U}(1)_-$, and show that it indeed has the same divergence as the field theory. We argue that this divergence has an infrared interpretation related to the closed string tadpoles generated by the D-branes. Therefore this divergence will be cancelled when the tadpoles are cancelled, which one must do in a consistent compact model.

In section 5, we will use $\mathcal{N} = 1$ string dualities to discuss the relation of this result to the heterotic string results of [20, 24].

3.1 Description of the model

We will study two D6-branes in type-IIA string theory, intersecting along $\mathbb{R}^4$. We write the $10d$ space as $\mathbb{R}^4 \times \mathbb{C}^3$: here $\mathbb{C}^3$ can be thought of as a local model for a nonsingular region of a Calabi-Yau manifold M. The D6-branes will fill out three-dimensional submanifolds of M. In a general Calabi-Yau threefold, $\mathcal{N} = 1$ supersymmetry requires that these submanifolds

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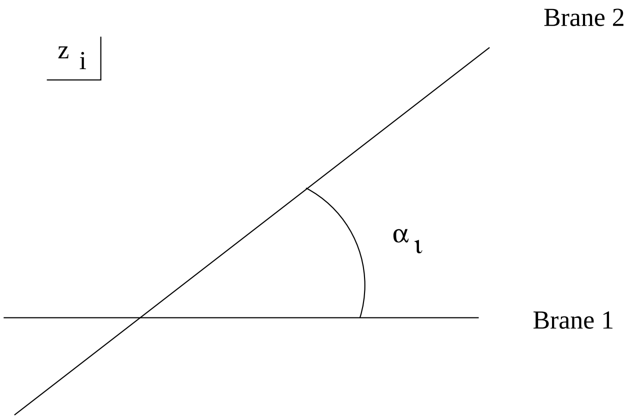

Figure 1: Two D6-branes intersecting along $\mathbb{R}^4 \subset \mathbb{R}^4 \times \mathbb{C}^3$, with one of the coordinates $z^{i=1,2,3}$ in $\mathbb{C}^3$ shown.

be special lagrangian [48]. The conditions for $\Sigma \subset M$ to be special lagrangian can be written in terms of the holomorphic $(3,0)$ form as:

$$\begin{aligned} \mathrm{Re}\, e^{i\theta}\Omega|_\Sigma &= \mathrm{vol}\ \Sigma \\ \mathrm{Im}\, e^{i\theta}\Omega|_\Sigma &= 0 \end{aligned} \tag{3.2}$$

where θ is some angle, often called the "phase" of Σ. In the present case, we can write a Kähler form on $\mathcal{I}^3$ as:

$$\omega = \eta_{i\bar{\jmath}} dz^i d\bar{z}^{\bar{\jmath}} \tag{3.3}$$

and a holomorphic $(3,0)$ form as:

$$\Omega = \epsilon_{ijk} dz^i dz^j dz^k\,. \tag{3.4}$$

Here $z^{i=1,2,3}$ are the canonical flat holomorphic coordinates on $\mathbb{C}^3$. It is easy to see that the cycles $\mathrm{Im}\, e^{i\theta^i} z^i = 0$ are special lagrangian cycles of $\mathbb{C}^3$ with phase $\theta = \sum_i \theta_i$. Two D6-branes wrapped on two intersecting cycles Σ_1, Σ_2 with the same phase will together preserve $d = 4, \mathcal{N} = 1$ supersymmetry as well, as the union $\Sigma_1 \bigcup \Sigma_2$ satisfies the above conditions.

In our local model, we will take $\Sigma_1 \in \mathbb{C}^3$ to be the submanifold $\mathrm{Im}\, z^i = 0$, and Σ_2 to be the submanifold $\mathrm{Im}\, e^{i\theta_i} z^i = 0$. Of course $\mathbb{C}^3$ has *vanishing* holonomy, so the theory has 32 supercharges instead of eight before the D-branes are added. We will label these branes "1" and "2" (see figure 1). Still, while Σ_1 by itself preserves $\mathcal{N} = 4$ SUSY in four dimensions, if we choose $\sum_i \theta_i = 0$, $\theta_i \neq 0\ \forall i$, then the branes Σ_1, Σ_2 together preserve $\mathcal{N} = 1$ spacetime SUSY in $\mathbb{R}^4$ [21]. The light spectrum of $4d$ fields was worked out in [21]. There is a $\mathrm{U}(1)^2 = \mathrm{U}(1)_1 \times \mathrm{U}(1)_2$ gauge symmetry, and a single chiral multiplet with charge $(1,-1)$. The off-diagonal combination $\mathrm{U}(1)_-$ of $\mathrm{U}(1)_1$ and $\mathrm{U}(1)_2$ is therefore anomalous. The anomaly is cancelled by anomaly inflow due to couplings of the D-brane to the RR potentials. The coupling $\int_{M_6} F \wedge C^{(5)}_{RR}$ for D6-branes on M_6 leads in particular to the coupling

$$\left(\int_{\mathbb{R}^4 \times \Sigma_1} + \int_{\mathbb{R}^4 \times \Sigma_2} \right) F \wedge C^{(5)}\,. \tag{3.5}$$

If F_1, F_2 are the gauge fields for branes 1 and 2, and we define the following two-forms in $\mathbb{R}^4$:

$$C_{1,2} = \int_{\Sigma_{1,2}} C^{(5)} \,,$$

then (3.5) leads to the couplings in the 4d effective action

$$S_{\text{anom}} = \frac{1}{2} \int d^4x \,[(F_1 - F_2) \wedge (C_1 - C_2) + (F_1 + F_2) \wedge (C_1 + C_2)] \,. \tag{3.6}$$

The first term on the right hand side leads to a 4d description of anomaly cancellation via the Green-Schwarz mechanism. It will become important in section 4.

3.2 Computing the one-loop contribution to the D-term.

The diagonal linear combination $\mathrm{U}(1)_+$ of $\mathrm{U}(1)_1 \times \mathrm{U}(1)_2$ should get no one-loop contribution to the corresponding FI D-term, as there are no charged fields are coupled to it perturbatively. On the other hand, according to the nonrenormalization theorem in [19], the FI D-term for the off-diagonal combination $\mathrm{U}(1)_-$ *will* get a one-loop contribution in the field theory limit. We will now test this via a direct computation in string theory.

In the RNS formalism, the FI D-term is proportional to the one-point function

$$\delta\xi = \sum_{a,b,i} \int dtdw \, \mathrm{tr}_{ab,i} \, \frac{1}{2} \left(1 + (-1)^F\right) V_{D,aa}(w) e^{-2\pi t L_0} (-1)^{F_{st}} \equiv \sum_{ab} A_{ab} \tag{3.7}$$

on the cylinder. The sum is over the Chan-Paton factors a, b of the two boundaries, and over the periodicity $i = (\mathrm{NS}, \mathrm{R})$ of the fermions. F_{st} is the spacetime fermion number — this leads to a factor of -1 in the Ramond sector — and F is the worldsheet fermion number. a denotes the Chan-Paton index of the boundary at $\sigma = 0$, and b the Chan-Paton index of the boundary at $\sigma = \pi$. The trace is over the oscillator modes and zero modes of the worldsheet fields. L_0 the zero mode of the worldsheet energy-momentum tensor, t is the modular parameter of the cylinder, w a location on the boundary of the cylinder, and $V_{D,ab}$ the vertex operator (2.1) for the FI D-term corresponding to $\mathrm{U}(1)_-$. For this calculation, one may either integrate V_D around the boundary of the cylinder, or divide by the length of this boundary and fix the position of V_D. Either way, one may write (3.7) as

$$\delta\xi = \mathrm{tr} \int dt \frac{1}{2} (1 + (-1)^F)(-1)^{F_{st}} V_{D,aa,0} e^{-2\pi t L_0} \,. \tag{3.8}$$

The vertex operator for the auxiliary field in the vector multiplet for $\mathrm{U}(1)_D$ is:

$$V_{D,aa',0} = J_0 \begin{pmatrix} 1 & 0 \\ 0 & 1 \end{pmatrix}_{aa'} \,, \tag{3.9}$$

while the vertex operator for the auxiliary field for $\mathrm{U}(1)_-$ is:

$$V_{D,aa',0} = J_0 \begin{pmatrix} 1 & 0 \\ 0 & -1 \end{pmatrix}_{aa'} \,. \tag{3.10}$$

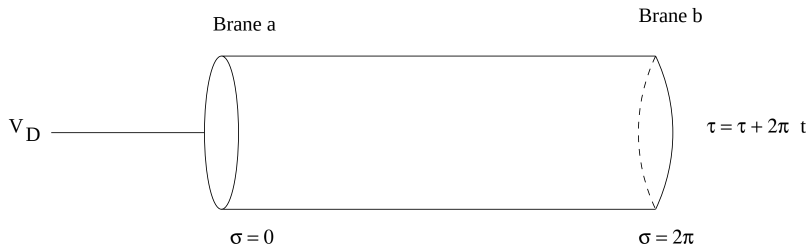

Figure 2: The one-loop diagram for V_D. Here $a, b = 1, 2$ labels whether each boundary ends on brane 1 or 2, respectively.

Here J_0 is the zero mode of the $\mathrm{U}(1)_R$ current, and the matrix acts on the Chan-Paton indices $a = 1, 2$ of the boundary at $\sigma = 0$ in which V_D is inserted. The indices $a = 1, 2$ denote branes 1 and 2, respectively. (See figure 2.)

We will find it useful to rewrite the amplitude (3.7) as

$$\begin{aligned} \langle V_D \rangle &= \frac{-i}{2\pi}\partial_\nu \sum_{abi} \mathrm{tr}_{ab,i} \int dt \frac{1}{2}(1+(-1)^F)(-1)^{F_{st}} s_{aa} e^{2\pi i \nu J_0} e^{-2\pi t L_0}|_{\nu=0} \\ &\equiv \frac{-i}{2\pi}\partial_\nu \, \mathrm{tr} \int dt \frac{1}{2}(1+(-1)^F)(-1)^{F_{st}} s_{aa} z^{J_0} q^{L_0}|_{\nu=0} \,, \end{aligned} \tag{3.11}$$

where

$$z \equiv e^{2\pi i \nu} \,; \qquad q \equiv e^{-2\pi t} \equiv e^{2\pi i \tau} \,.$$

We can break up the amplitude into Chan-Paton sectors:

$$\begin{aligned} \delta\xi_+ &= \frac{-i}{2\pi}\partial_\nu \int dt \, (A_{11}(t) + A_{12}(t) + A_{21}(t) + A_{22}(t)) \\ \delta\xi_- &= \frac{-i}{2\pi}\partial_\nu \int dt \, (A_{11}(t) + A_{12}(t) - A_{21}(t) - A_{22}(t)) \,, \end{aligned} \tag{3.12}$$

where ξ_+, ξ_- are the FI D-terms for $\mathrm{U}(1)_+$ and $\mathrm{U}(1)_-$, respectively. We will now discuss each of the A_{ab}s in turn.

A_{11} and A_{22}. Strings in the 11 and 22 sectors are neutral with respect to both $\mathrm{U}(1)_1$ and $\mathrm{U}(1)_2$. Therefore, we expect them to give a vanishing contribution to the FI D-terms. We can see this by simply examining the trace over fermion modes in the $\mathbb{C}^3$ direction. The point is that these fermions are the only fields on the worldsheet charged under the worldsheet $\mathrm{U}(1)_R$, and so lead to all of the ν-dependence of the amplitudes. For the 11 sector, we can use the mode expansions described in appendix A to find:

$$\begin{aligned} \mathrm{tr}_{11,NS}(-1)^{F_{st}} z^{J_0} q^{L_0} &= \prod_{n=1}^{\infty} (1 + zq^{n-\frac{1}{2}})^3 (1 + z^{-1}q^{n-\frac{1}{2}})^3 \\ \mathrm{tr}_{11,NS}(-1)^{F+F_{st}} z^{J_0} q^{L_0} &= \prod_{n=1}^{\infty} (1 - zq^{n-\frac{1}{2}})^3 (1 - z^{-1}q^{n-\frac{1}{2}})^3 \\ \mathrm{tr}_{11,R}(-1)^{F_{st}} z^{J_0} q^{L_0} &= -(z^{1/2} + z^{-1/2}) \prod_{n=1}^{\infty} (1 + zq^n)^3 (1 + z^{-1}q^n)^3 \,. \end{aligned} \tag{3.13}$$

The expressions for the 22 sector are identical. We do not write out $\mathrm{tr}_{11,R}(-1)^{F+F_{st}}(\cdots)$ because the spacetime fermion zero modes lead to a vanishing contribution from this sector.

All of the expressions in (3.13) are even under $\nu \to -\nu$, and regular at $\nu = 0$. Therefore, the derivative of any of these expressions with respect to ν must vanish at $\nu = 0$, and so $A_{11} = A_{22} = 0$.

A_{12} **and** A_{21}. The strings stretching between branes 1 and 2 are charged under $U(1)_-$. We expect any contributions to $\delta\xi_-$ to come from the 12 and 12 Chan-Paton sectors. We will work through the computation of A_{12} in some detail. It will be clear from the form of the answer that $\partial_\nu A_{21} = -\partial_\nu A_{12}$.

Let us begin with the trace over the spacetime fields and ghosts. The trace over the superconformal ghosts will cancel the trace over the fermion modes and bosonic oscillator modes for the "longitudinal" fields X^0, X^1, ψ^0, ψ^1 (see appendix A for notation). The contribution to A_{12} from the spacetime bosons and the bosonic ghosts is:[4]

$$A_{12,\mathrm{s.t.bosons}} = \int \frac{d^4p}{(2\pi)^4} e^{-\pi t p^2/2} \frac{q^{\frac{1}{12}}}{\eta(\tau)^2} q^{E_{gs,1}} = \left(\frac{1}{\pi^2 t}\right)^2 \frac{q^{\frac{1}{12}}}{\eta(\tau)^2}\,, \tag{3.14}$$

where $E_{gs,1}$ is the vacuum energy of these fields. The nonzero contributions of the spacetime fermions and fermionic ghosts are:

$$\begin{aligned}
\mathrm{tr}_{12,NS,\mathrm{s.t.fermions}}\, z^{J_0} q^{L_0} &= q^{\frac{1}{24}+E_{gs,3,NS}} \frac{\vartheta_{00}(0,\tau)}{\eta(\tau)} \\
\mathrm{tr}_{12,NS,\mathrm{s.t.fermions}} (-1)^F z^{J_0} q^{L_0} &= q^{\frac{1}{24}+E_{gs,3,NS}} \frac{\vartheta_{01}(0,\tau)}{\eta(\tau)} \gamma_3 \\
\mathrm{tr}_{12,R,\mathrm{s.t.fermions}}\, z^{J_0} q^{L_0} &= -q^{-\frac{1}{12}+E_{gs,3,R}} \frac{\vartheta_{10}(0,\tau)}{\eta(\tau)}\,,
\end{aligned} \tag{3.15}$$

where $E_{gs,3,NS}$ and $E_{gs,3,R}$ are the ground state energies of these modes in the Neveu-Schwarz and Ramond sectors, respectively, and $(-1)^F|0,NS\rangle = \gamma_3|0,NS\rangle$ in the NS sector.

The contribution of the "internal" complex worldsheet bosonic fields $Z^{i=1,2,3}$ is:

$$A_{12,\mathrm{int.bosons}} = (-i)^3 q^{\frac{1}{4}} \frac{\eta(\tau)^3}{\prod_{i=1}^3 \vartheta_{11}(|\alpha_i|\tau,\tau)} q^{E_{gs,2}}\,, \tag{3.16}$$

where $E_{gs,2}$ is the vacuum energy of these fields. As in appendix A, $\alpha_i = \frac{\theta_i}{2\pi}$, where θ_i is the angle between branes 1 and 2 in the z^i-plane of $\mathbb{C}^3$.

Finally, the contribution of the "internal" complex fermions is:

$$\begin{aligned}
\mathrm{tr}_{12,NS,\mathrm{int.fermions}} (-1)^{F_{st}} z^{J_0} q^{L_0} &= q^{\frac{1}{8}+E_{gs,4,NS}} \frac{\prod_{i=1}^3 \vartheta_{00}(\nu+\alpha_i\tau,\tau)}{\eta(\tau)^3} \\
\mathrm{tr}_{12,NS,\mathrm{int.fermions}} (-1)^{F+F_{st}} z^{J_0} q^{L_0} &= q^{\frac{1}{8}+E_{gs,4,NS}} \frac{\prod_{i=1}^3 \vartheta_{01}(\nu+\alpha_i\tau,\tau)}{\eta(\tau)^3} \gamma_4 \\
\mathrm{tr}_{12,R,\mathrm{int.fermions}} (-1)^{F_{st}} z^{J_0} q^{L_0} &= q^{-\frac{1}{4}+E_{gs,4,R}} \frac{\prod_{i=1}^3 \vartheta_{10}(\nu+\alpha_i\tau,\tau)}{\eta(\tau)^3}\,.
\end{aligned} \tag{3.17}$$

[4]We use the conventions in [46] for $\eta(q)$ and for the theta functions ϑ_{ab}.

Again, $E_{gs,4,NS}$ and $E_{gs,4,R}$ are the ground state energies in the NS and R sectors, respectively, and $(-1)^F|0,NS\rangle = \gamma_4|0,NS\rangle$.

Putting these all together, we find that:

$$A_{12} = \left(\frac{1}{\pi^2 t}\right)^2 \frac{(-i)^3}{\eta(\tau)^3 \prod_{i=1}^3 \vartheta_{11}(\alpha_i\tau,\tau)} \times \\ \times \left[\vartheta_{00}(0,\tau)\prod_{i=1}^3 \vartheta_{00}(\alpha_i\tau,\tau) - \vartheta_{01}(0,\tau)\prod_{i=1}^3 \vartheta_{00}(\alpha_i\tau,\tau) - \vartheta_{10}(0,\tau)\prod_{i=1}^3 \vartheta_{00}(|\alpha_i|\tau,\tau)\right] \tag{3.18}$$

Here we have used the fact that $(-1)^F|0,NS\rangle = \gamma_3\gamma_4|0,NS\rangle = -|0,NS\rangle$, that $\sum_{k=1}^4 E_{gs,k,NS} = -1/2$, and that $\sum_{k=1}^4 E_{gs,k,R} = 0$.

This expression can be simplified via the Riemann theta identities (q.v. eq. (13.4.20-21) of [46] or section 1, section 5, of [49]) to:

$$A_{12} = \left(\frac{1}{\pi^2 t}\right)^2 \frac{(-i)^3\vartheta_{11}(\frac{3\nu}{2},\tau)\prod_{i=1}^3 \vartheta_{11}(-\frac{\nu}{2}+\alpha_i\tau,\tau)}{\eta(\tau)^3\prod_{i=1}^3 \vartheta_{11}(|\alpha_i|\tau,\tau)}\,. \tag{3.19}$$

Next, we wish to compute $\partial_\nu A_{12}|_{\nu=0}$. Because $\vartheta_{11}(0,\tau) = 0$,

$$\frac{-i}{2\pi}A_{12}|_{\nu=0} = \left(\frac{1}{\pi^2 t}\right)^2 \frac{(\partial_\nu\vartheta_{11}(\frac{3\nu}{2},\tau))_{\nu=0}\prod_{i=1}^3 \vartheta_{11}(\alpha_i\tau,\tau)}{2\pi\eta(\tau)^3\prod_{i=1}^3 \vartheta_{11}(\alpha_i\tau,\tau)} = \frac{18}{\pi t^2}\,. \tag{3.20}$$

where we have used $\partial_\nu\vartheta_{11}(\nu,\tau)|_{\nu=0} = (-2\pi\eta(\tau))^3$. We have also used the fact that one of the angles α is negative, which we have chosen to be $\alpha_3 < 0$: in this case,

$$\frac{\vartheta_{11}(\alpha_3\tau,\tau)}{\vartheta_{11}(|\alpha_3|\tau,\tau)} = -1\,,$$

which contributes an additional minus sign, leading to the overall sign in (3.20).

Inspection of (3.17) reveals that A_{12} is invariant under the combined operation $\alpha_i \to -\alpha_i$, $\nu \to -\nu$. In the mode expansion for the 21 sector, the only difference from the 12 sector is that the angles have the opposite sign. One may therefore write $A_{21}(\nu,t) = A_{12}(-\nu,t)$, and so

$$\partial_\nu A_{21}|_{\nu=0} = -\partial_\nu A_{12}|_{\nu=0}\,. \tag{3.21}$$

3.3 Physical interpretation of the one-loop amplitudes for $\delta\xi_{+,-}$

Using eq. (3.20) and (3.21), we find:

$$\begin{aligned} \delta\xi_+ &= \frac{-i}{2\pi}\partial_\nu\int dt\,(A_{12}+A_{21}) = 0 \\ \delta\xi_- &= \frac{-i}{2\pi}\partial_\nu\int dt\,(A_{12}-A_{21}) = \frac{36}{\pi}\int_\epsilon^\infty dt\frac{1}{t^2}\,, \end{aligned} \tag{3.22}$$

where $\epsilon = \Lambda_{UV}^{-2}$ is the open-string-channel ultraviolet cutoff. These answers are consistent with the results of [19]. As expected, there is no correction to the FI D-term for U(1)$_+$, since there is no massless chiral multiplet charged under this U(1). For $\delta\xi_-$, the form

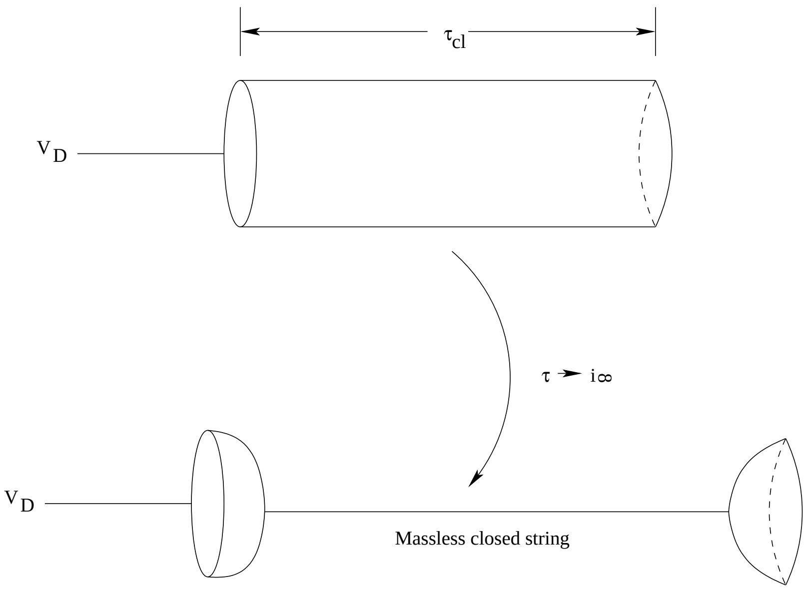

Figure 3: Factorization of the one-loop correction to ξ_-.

of (3.20) indicates that none of the oscillator modes contribute. The fact that the oscillator contributions have cancelled reflects the fact that massive fields do not renormalize the FI term; this is an index quantity. The zero modes for the massless chiral multiplet charged under $U(1)_-$ lead to a quadratic ultraviolet divergence.

Unlike the heterotic string, this divergence is not cut off by any modular group action. Such a divergence must have an infrared interpretation. Indeed, if one performs a modular transformation $\tau_{cl} = -1/\tau = i/t$ to the closed string channel (see figure 3), then we find that

$$\delta\xi_- = \frac{-36i}{\pi} \int_0^{i\epsilon} d\tau_{cl}\,. \tag{3.23}$$

This represents an infrared divergence in the closed string channel. This divergence is due to the exchange of massless closed string modes in the factorization limit illustrated in figure 3. The divergence from A_{12} (A_{21}) is proportional to the tadpoles generated by D6-brane 2 (D6-brane 1). It will therefore be cancelled when the tadpoles are cancelled.[5] This is consistent with the result of [25].

A remaining question is whether any finite correction to ξ_- can remain. If the D-brane tadpoles are cancelled by perturbative orientifolds, the correction will vanish exactly, as in [25], if the orientifold projection preserves SUSY. The nonrenormalization theorem of [19], and the form of the string amplitude (3.20), imply that the contribution of each massive supermultiplet vanishes separately. This will remain true for fields which survive the orientifold projection. The orientifold sectors will give no contributions to ξ_- from open string oscillator modes, and the divergence in these sectors has the same form $c\int dt/t^2$ for

[5]In noncompact examples, this can be accomplished by solving the equations of motion for the RR 7-form potential in the presence of the D6-brane source. We thank Atish Dabholkar and Howard Schnitzer for bringing this fact to our attention.

some constant c. If the orientifold cancels the divergence, there is no room for a finite part. This is a mystery from the standpoint of the low-energy effective field theory, unless there is a symmetry which would forbid such a term. It is tempting to blame the underlying $\mathcal{N} = 2$ structure of the closed-string physics; we leave this question for future work.

4. Cosmic D-term strings are D-branes

The local model we studied in the previous example is one where the FI D-term is vanishing at tree level. The local model in type-IIA for an open string system with a single charged chiral scalar field, and tree-level FI D-term near zero, has been described mathematically by [30] and given the physical interpretation we will use in [29] (see also [50] for a nice review). As in the previous section, the low-energy gauge group of the model is $\mathrm{U}(1)_1 \times \mathrm{U}(1)_2 = \mathrm{U}(1)_+ \times \mathrm{U}(1)_-$. There is a chiral superfield Φ with scalar component ϕ, charged under $\mathrm{U}(1)_-$. The D-term potential is:

$$V_D = e^2 \left(|\phi|^2 - \xi\right)^2 . \tag{4.1}$$

When $\xi < 0$, $V_D \neq 0$ and SUSY is broken; meanwhile, at the minimum of V_D, $\phi = 0$ and both $\mathrm{U}(1)_+$ and $\mathrm{U}(1)_-$ are unbroken. When $\xi > 0$, there is an S^1 vacuum manifold at $|\phi|^2 = \xi$. SUSY is unbroken, but $\mathrm{U}(1)_-$ is Higgsed in the vacuum.

This setup can be realized by D6-branes in type-IIA, as described by [30, 29]. The D6-branes wrap various special lagrangian three-cycles $\Sigma \subset M_6$, where M_6 is a local model of a Calabi-Yau background. The geometric description of the system is show on the left in figure 4. $\mathrm{U}(1)_1$ is the gauge group for a D6-brane wrapped on the three-cycle N_+ and $\mathrm{U}(1)_2$ is the gauge group for a D6-brane on cycle N_-. ξ is determined by complex structure moduli. For $\xi < 0$, the cycles $N_\pm$ have different phases in the sense of (3.2). As $\xi \to 0$, the strings stretching between N_+ and N_- include a light chiral multiplet charged under $\mathrm{U}(1)_-$. At $\xi = 0$, this multiplet becomes massless, and SUSY is restored: the phases of

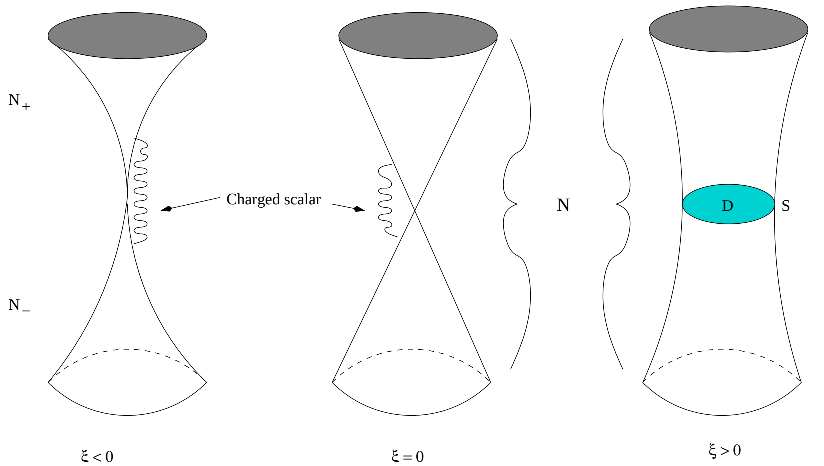

Figure 4: The local geometry of the intersection in the $\xi < 0$ broken-SUSY phase, the $\xi = 0$ phase boundary, and the $\xi > 0$ Higgs phase.

$N_{\pm}$ become identical. When $\xi > 0$, the lightest chiral scalar becomes tachyonic. If one condenses this scalar field, the two three-cycles $N_{\pm}$ merge into the three-cycle N which is equivalent in homology to $N_+ + N_-$.

In the $\xi > 0$ phase, there should be cosmic string solutions, which we will call "D-term strings". Dvali and collaborators [31, 32] have shown that such solutions have the following properties.

1. The D-term strings have tension proportional $T = 2\pi\xi$.

2. The D-strings are BPS, with chiral $\mathcal{N} = (2,0)$ worldsheet supersymmetry.

3. The D-term string should be a magnetic flux tube for the anomalous U(1) under which ϕ is charged. If this anomaly is cancelled via coupling to an axion a:

$$A^{\mu}_{-}\partial_{\mu}a \tag{4.2}$$

   with $a \sim a + 2\pi$, then a should shift by 2π around an S^1 which circles the D-term string once.

Dvali et al. [31] argued that the cosmic strings which can appear in D-$\bar{D}$ brane annihilation, such as in "brane inflation" scenarios [9]–[16], are in fact D-term strings. More generally, they showed that the inflaton potential energy in such scenarios has a D-term component. These are potentially advantageous models. SUSY is broken during inflation, and if F-terms contribute significantly to the inflaton potential, this breaking makes the inflaton dangerously heavy unless the theory is somewhat finely tuned [51, 52]. If the inflaton potential arises from an FI D-term and F-terms do not contribute significantly, no fine-tuning is required in order to keep the inflaton potential flat [17, 18].

In the case that D-term inflation arises from brane-antibrane annihilation, the D-term ξ scales inversely with the string coupling itself, and vanishes only at strong coupling. For D6-branes in type-IIA Calabi-Yau backgrounds (the only branes filling out $\mathbb{R}^4$ which potentially preserve $\mathcal{N} = 1, d = 4$ SUSY), the scenario described above is more general, in the sense that the FI coupling can be a function of all of the hypermultiplet moduli. The basic issue, as systematized by Douglas [7], is that for D-branes in Calabi-Yau backgrounds, the notion of "brane" and "antibrane" depends on the closed string fields descending from $\mathcal{N} = 2$ hypermultiplets — complex structure moduli, in this case. Therefore, if one can stabilize these moduli along the lines of [3] with some control over their vevs, one has great freedom in designing D-term potentials.[6] As a brane inflation scenario, this is the embedding into a Calabi-Yau background of the scenario in [54, 14].

A natural question is, then, the identity of the D-term strings in these models. A puzzle is that a natural way to get a light cosmic string would be by wrapping a $(p+1)$-brane around a vanishing p-cycle $\Sigma_p \subset M_6$. But the transitions described in [30, 29, 7] generically occur in perfectly regular interior points of the closed-string moduli space.

[6]In general the F-terms introduced by stablization of the Kähler moduli can overwhelm the D-terms [53, 32]; one must take care in constructing actual models.

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Inspired by [35, 36], we identify the D-term strings for the model described in figure 4 in the following way. As $\xi \to 0_+$, the SUSY 3-cycle develops a pinch; the local geometry near the pinch is a "Lawlor neck" (q.v. [30, 50]), topologically $S^2 \times \mathbb{R}$. Each S^2 bounds a 3-ball. There is a minimal-volume, special lagrangian 3-ball, called D in figure 4, which is bounded by the smallest S^2 $S = \partial D$ in the neck. D goes to zero volume as $\xi \to 0$, and it has phase i, in the sense of (3.2), relative to N. Now, a D4-brane can consistently end on a D6-brane along a submanifold of codimension 3. Therefore, a D4-brane with worldvolume $D \times \mathbb{R}^{1,1} \subset M_6 \times \mathbb{R}^4$ is a candidate for the D-term string, as it has vanishing tension in the four-dimensional theory when $\xi \to 0, \mathrm{vol}(D) \to 0$. This is consistent with property (1) above. In the remainder of this section we will show that this string has properties (2), (3) listed above, and so is a good candidate for the D-term string in the model of [30, 29].

4.1 Worldsheet SUSY for the D-string

Choose the overall phase of the holomorphic three-form Ω so that N has phase 0, i.e. $\mathrm{Im}\,\Omega|_N = 0$. Let Σ be the spectral flow operator on for the open string which generates 1/2-unit of spectral flow from $NS \to R$. The operator generating a full unit of spectral flow $NS \to NS$ is

$$\Sigma^2 = \Omega_{ijk}\psi^i\psi^j\psi^k\,. \tag{4.3}$$

The boundary conditions on ψ, Ω then imply that $\Sigma_+ = \bar{\Sigma}_-$ where $+,-$ denote left- and right-movers as in appendix A. This boundary condition is consistent with some general facts regarding the $\mathcal{N} = 2$ worldsheet CFT. The $\mathrm{U}(1)_R$ current can be written in terms of a worldsheet boson, $J_{R,\pm} = c\partial_\pm H$. We can then write $\Sigma_\pm = e^{iaH_\pm}$ For the A-type boundary conditions required for the D6-branes to preserve $\mathcal{N} = 1$ SUSY in $d = 4$ [48, 55], $J_{R,+} = -J_{R,-}$.

Next, let $S_{\alpha,\pm}$ be the spin fields for the spacetime fermions, so that the currents implementing $\mathcal{N} = 2$ spacetime SUSY for the closed string theory are:

$$\begin{aligned} Q_{\alpha,+} &= e^{-\phi/2} S_{\alpha,+}\Sigma_+ \\ Q_{\alpha,-} &= e^{-\tilde{\phi}/2} S_{\alpha,-}\bar{\Sigma}_-\,. \end{aligned} \tag{4.4}$$

If all directions in spacetime are Neumann, then one can show that $S_{\alpha,+} = S_{\alpha,-}$. The boundary conditions on $\Sigma_\pm$ then imply that

$$Q_\alpha = e^{-\phi/2}\Sigma_+ S_{\alpha,+} - e^{-\tilde{\phi}/2}\bar{\Sigma}_- S_{\alpha,-} \tag{4.5}$$

and the corresponding operator for $\bar{Q}_{\dot{\alpha}}$ generate the $\mathcal{N} = 1$ spacetime supersymmetry preserved by the D6-branes.

We would now like to understand which of these SUSYs are preserved in the presence of our candidate D-term string. Therefore, $\Sigma_L = -i\bar{\Sigma}_R$ for strings ending on this D4-brane. Q_α in (4.5) is then preserved if $S_{\alpha,L} = iS_{\alpha,R}$. This is in fact a consequence of worldsheet SUSY and the boundary conditions on the spacetime coordinates for strings ending on the D4-brane. In spacetime, one can write the spacetime spin fields as:

$$\begin{aligned} S_{1,\pm} &= e^{i(H_{1,\pm}+H_{2,\pm})/2}\,, \qquad & S_{2,\pm} &= e^{-i(H_{1,\pm}+H_{2,\pm})/2} \\ S_{\dot{1},\pm} &= e^{i(H_{1,\pm}-H_{2,\pm})/2}\,, \qquad & S_{\dot{2},\pm} &= e^{-i(H_{1,\pm}-H_{2,\pm})/2} \end{aligned} \tag{4.6}$$

where

$$\begin{aligned} e^{iH_1} &= \psi^t + i\psi^z \\ e^{iH_2} &= \psi^2 = \psi^x + i\psi^y\,. \end{aligned} \tag{4.7}$$

Here t, x, y, z are coordinates on $\mathbb{R}^4$. Assume that the string is stretching along z. ψ^t, ψ^z are the worldsheet superpartners of scalars with Neumann boundary conditions, and ψ^x, ψ^y are superpartners of a scalars with Dirichlet boundary conditions, so

$$\psi_+^{t,z} = \psi_-^{t,z}\,; \qquad \psi_+^{x,y} = -\psi_-^{x,y}\,. \tag{4.8}$$

This boundary condition is then consistent with the boundary condition $e^{iH_{1,+}} = e^{iH_{1,-}}$, $e^{iH_{2,+}} = ie^{iH_{2,-}}$. In terms of the spin fields, this implies:

$$\begin{aligned} S_{1,+} &= iS_{1,-}\,; \qquad S_{2,+} = -iS_{2,-} \\ S_{\dot{1},+} &= -iS_{\dot{1},-}\,; \qquad S_{\dot{2},+} = iS_{\dot{2},-}\,. \end{aligned} \tag{4.9}$$

This means that of the four supercharges in (4.5) and its conjugate, $Q_1, Q_{\dot{2}}$ are preserved. These supersymmetries are both right-moving along z. Therefore the D4-brane on $D \times \mathbb{R}^{1,1}$ has $\mathcal{N} = (2,0)$ SUSY in $\mathbb{R}^{1,1}$, consistent with our identification of this D4-brane as the D-term string.

4.2 Magnetic flux and axion charge carried by the D-string

As stated above, for the D4-brane on $D \times \mathbb{R}^{1,1}$ to be the D-term string we claim it is, it should carry magnetic flux under the anomalous gauge field; in other words, $\int_{\mathbb{R}^2} F_- = 1$, where $\mathbb{R}^2$ is the plane transverse to the string in four dimensions. When the anomaly in $\mathrm{U}(1)_-$ is cancelled via the Green-Schwarz mechanism through the coupling (4.2), this is gauge-equivalent to the statement that the string should have axion charge. This axion charge can be seen as follows. The specific coupling which takes the form (4.2) is the dimensional reduction of the Wess-Zumino coupling [23]

$$S_{7,GS} = \sum_i \int_{N_i \times \mathbb{R}^4} F_i \wedge C^{(5)} \tag{4.10}$$

where F_i is the worldvolume gauge field strength on the ith D6-brane, $C^{(5)}$ the 5-form RR potential of type-IIA string theory, and $\{N_i\}$ is the collection of D6-branes. If the D6-branes wrap $N = N_+ + N_-$, and for modes of F which are independent of the internal space, and are polarized along $\mathbb{R}^4$, (4.10) can be written as[7]

$$\begin{aligned} S_{4,GS} &= \frac{1}{2}\int d^4x \left[(F_+ - F_-)\wedge(C_+^{(2)} - C_-^{(2)}) + (F_+ + F_-)\wedge(C_+^{(2)} + C_-^{(2)})\right] \\ &= \frac{1}{2}\int d^4x \left(A_-^\mu \partial_\mu a_- + A_+^\mu \partial_\mu a_+\right). \end{aligned} \tag{4.13}$$

[7] This is completed to the object

$$L \ni (\partial a_- + A_-)^2 \tag{4.11}$$

which is invariant under the gauge transformation

$$a \mapsto a - \lambda, \quad A_\mu \mapsto A_\mu - \partial_\mu \lambda\,. \tag{4.12}$$

In unitary gauge, $a_- = 0$, and (4.11) is a mass term for the relative gauge field A.

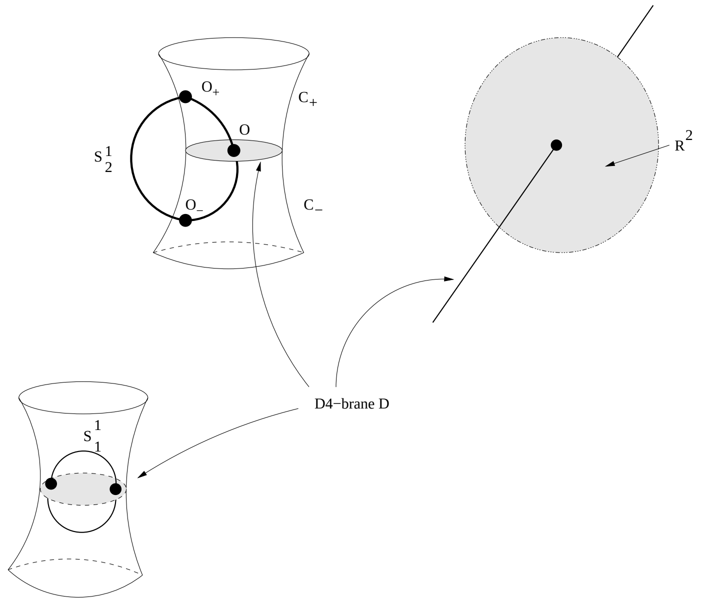

Figure 5: The ingredients involved in the construction of $\Sigma_{\mathrm{in,out}}$, used to deduce the magnetic flux and axion charge carried by the string.

Here $C^{(2)}_{\pm} = \int_{N_{1,2}} C^{(5)}$, and in four dimensions, $da_{\pm} = \star_4 d(C^{(2)}_{+} \pm C^{(2)}_{-})$. If the D4-brane on $D \times \mathbb{R}^{1,1}$ is the D-term string we claim it is, the worldvolume $\mathbb{R}^{1,1}$ should couple to $C^{(2)}_{+} - C^{(2)}_{-}$, so that a_{-} should shift by 2π around any S^1 encircling the string in $\mathbb{R}^4$. Because of (4.13), this will be true if the string carries magnetic flux under $\mathrm{U}(1)_{-}$.

We can find the magnetic flux carried by our D-string by adapting the arguments in [56]. The geometric setup is shown in figure 5. One must imagine an extra S^2 at each point in the figure, so that D is a three-ball, and its boundary a 2-sphere in N. Finally, there is an additional two-manifold $K \subset \mathbb{C}^3$ which is transverse to both N, D, and which is in general noncompact. Pick a point $P \in K$ lying on D.

We can construct a 5-cycle Σ by gluing the following 5-chains along their boundaries:

1. Σ_{in}. Take the spacelike disc $B_1 \subset \mathbb{R}^{3,1}$, which intersects the string once in spacetime at a point Q, and $K_1 \subset K$ which is a disc with P in the interior. Finally, imagine a family γ_t of circles which interpolates from S^1_2 at $t = 0$ to S^1_1 at $t = 1$. Define Σ_{in} as an S^1 fibration of $B_1 \times K_1$, such that the S^1 is γ_0 at $(Q, P) \in B_1 \times K_1$, and γ_1 at the boundary $\partial(B_1 \times K_1) = \partial B_1 \times K_1 \cup B_1 \times \partial K_1$. As indicated in figure 5, this has positive intersection with the D4-brane worldvolume, Δ at (P, Q, O), positive intersection with C_+ at (P, Q, O_+), and negative intersection with C_- at (P, Q, O_-). It is called Σ_{in} because the D4-brane passes through it.

2. Σ_{out}. Now take the disc B_2 which intersects the string at Q', with the boundary $\partial B_2 = -\partial B_1$, so that B_1 and B_2 can be glued together to form a 2-sphere. Similarly, take K_2 to be a disc with $\partial K_2 = -\partial K_1$. Let $\Sigma_{\mathrm{out}} = B_2 \times K_2 \times \gamma_1$. This intersects Δ twice, but each time with opposite sign, so the total intersection is zero.

Now, $\Sigma \equiv \Sigma_{\rm in} \cup \Sigma_{\rm out}$, defined by gluing $\Sigma_{\rm in,out}$ along their boundaries in the obvious way, therefore has the same total intersection with Δ, N_+, and N_- as $\Sigma_{\rm in}$. Σ is a closed five-cycle, so that the integral $\int d\star dC^{(5)}$ must vanish. The equation of motion for $C^{(5)}$ is:[8]

$$d\star dC^{(5)} = (2\pi)^3 g_s \ell_s^3 \left(\sum_k \delta^{(5)}(\Sigma_k^{D4}) + \sum_i \delta^{(3)}(\Sigma_i^{D6}) \wedge F \right), \tag{4.14}$$

where Σ_k^{D4} are the set of cycles about which D4-branes are wrapped, and Σ_i^{D6} the set of cycles about which D6-branes are wrapped. Integrating this equation over Σ, the result for our configuration is:

$$\begin{aligned} 0 &= \int_\Sigma \left(\delta^{(5)}(\Sigma^{D4}) + \delta^{(3)}(N_+) \wedge F_+ + \delta^{(3)}(N_-) \wedge F_- \right) \\ &= \Sigma \cap \Delta + \left((K \times S^1_2) \cap N_+ \right) \int_{\mathbb{R}^2} F_+ + \left((K \times S^1_2) \cap N_- \right) \int_{\mathbb{R}^2} F_- \\ &= 1 + \int_{\mathbb{R}^2} (F_+ - F_-) \,. \end{aligned} \tag{4.15}$$

Therefore $\int_{\mathbb{R}^2}(F_+ - F_-) = -1$, and the string carries the magnetic flux consistent with our identification of this as a D-term string.

4.2.1 Axion charge

Next, we verify that we have identified the correct axion charge of the string. Recall that gauge flux F on the worldvolume of D6 carries D4 charge. Because of this, there is a sense in which the D4 doesn't actually end – the locus on which $C^{(4)}$ is sourced does not have a boundary – but rather, half of it goes into C_+ and (minus) half of it goes into C_-. In this sense, there is, in the higgs phase, actually a D4 wrapping the *closed* cycle $\Sigma \equiv \frac{1}{2}[N_+ - N_-] = \frac{1}{2}[C_+ - C_- + 2D]$, which is intersection-dual to $[N] = [N_+ + N_-]$. In equations, this statement can be tested by looking again at the equations of motion for the RR 5-form (4.14). Consider integrating this equation relating 5-forms over the 5-chain $\Sigma \times B_2$ where B_2 is a gaussian volume surrounding the string in spacetime (i.e. B_2 is a 2-ball surrounding the string in spacetime, whose boundary is an S^1 at infinity encircling the string). This gives

$$\int_{B_2} d\star d \int_\Sigma C^{(5)} = (2\pi)^3 g_s \ell_s^3 \int_{B_2} \delta^{(2)}(\text{string}) \int_\Sigma \delta^{(3)}(D) + (2\pi)^3 g_s \ell_s^3 \int_{B_2} F \int_\Sigma \delta^{(3)}(N) \,. \tag{4.16}$$

The left-hand side of this equation is $(2\pi\ell_s)^3 g_s$ because $\star d \int_\Sigma C^{(5)} = da$ is the axion flux, and $\int_{B_2} d(da_-) = \int_{S^1_\infty} da_- = (2\pi\ell_s)^3 g_s$: the axion goes once through its period when you go around the string.[9] The integral on the right-hand side of (4.16) is one because it is equal to the intersection product $\Sigma \cap N = 1$. Indeed, this is consistent with our claim that the D-term string is charged under the RR axion $a_- = \int_{N_+ + N_-} C^{(3)}_{RR}$.

[8] We follow the notation in, for example, ref. [57] in deducing factors of g, ℓ_s and 2π. The only difference is that we rescale $F \to (\sqrt{2\pi})\ell_s^2 F$, so that F is dimension 2 and has a standard kinetic term.

[9] The normalization follows by demanding that equation (4.14) is consistent in flat space in the presence of a flat D4-brane.

5. String duality and D-term strings

We conclude by discussing the relation of our results to the analogous physics of the heterotic string compactified on a Calabi-Yau threefold. In that theory, when the anomaly in a U(1) factor of the gauge group is cancelled by the Green-Schwarz mechanism involving the NS 2-form B, the corresponding D-term gets a one-loop correction proportional to the string tension M_s^2 [20, 47, 24]. It is worth noting that this is consistent with results about D-term strings in [31, 32]. A D-term of order M_s^2, if SUSY remains unbroken, implies a cosmic string with $N = (2,0)$ worldsheet supersymmetry, and tension M_s^2, whose worldsheet couples linearly to B. But this is the heterotic string itself! This is in keeping with the fact that the heterotic string can be written as a (singular) soliton of the massless fields of string theory [58].

It is also consistent with our results via string duality. Start with the duality between heterotic string theory compactified on T^3 and M-theory compactified on $K3$. The heterotic string is dual to an M5-brane wrapped on the K3. One may fiber these dual pairs over a (large) three-manifold base to find an $\mathcal{N} = 1$ dual pair of heterotic string theory compactified on a (T^3-fibered) Calabi-Yau background and M-theory compactified on a ($K3$-fibered) G_2 manifold Y. This M theory compactification can be reduced to IIA along an S^1-fibration of the K3 fibers of Y, leading to a CY backround of IIA with orientifold 6-planes and D6-branes at the loci where the S^1 fiber shrinks. The M5-brane wrapped on a K3 fiber becomes an open D4-brane ending on the D6-branes, as in section four. This string is then a D-term string for one of the D6-brane gauge groups, and has tension of order $1/g_s$, consistent with the fact that the D-term arises at type-IIA tree-level.

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A. Mode decomposition for open strings ending on intersecting D6-branes

This appendix contains the explicit mode decompositions for open strings ending on the D6-branes discussed in section 3. The results for worldsheet fields corresponding to the $\mathbb{R}^4$

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directions are standard — the fields are bosons satisfying Neumann boundary conditions, and their fermionic superpartners. We will concentrate on the internal bosons and fermions. The basic results are in [21]; we write them explicitly here to establish our conventions.

A.1 Internal bosons — closed string sector

The closed string modes along $\mathbb{C}^3$ are complex scalars

$$Z^i = z_0^i + p^i\tau + \frac{i}{\sqrt{4\pi}}\sum_{n=-\infty}^{\infty}\left(\frac{Z_n^i}{n}e^{-in(\tau-\sigma)} + \frac{\tilde{Z}_n^i}{n}e^{-in(\tau+\sigma)}\right) \tag{A.1}$$

where τ runs along the time axis of the cylinder, $\sigma \in [0, 2\pi]$, $i, \bar{i} = 1, 2, 3$ are complex indices, and the metric is the flat metric $\eta_{i\bar{j}}$ on $\mathbb{C}^3$. We will write the complex conjugate fields as

$$\bar{Z}^{\bar{i}} = \bar{z}_0^{\bar{i}} + \bar{p}^{\bar{i}}\tau + \frac{i}{\sqrt{4\pi}}\sum_{n=-\infty}^{\infty}\left(\frac{\bar{Z}_n^{\bar{i}}}{n}e^{-in(\tau-\sigma)} + \frac{\tilde{\bar{Z}}_n^{\bar{i}}}{n}e^{-in(\tau+\sigma)}\right). \tag{A.2}$$

The canonical commutation relations imply that

$$[Z_n^i, \bar{Z}_{n'}^{\bar{j}}] = n\delta_{n+n'}\eta^{i\bar{j}}\,; \qquad [\tilde{Z}_n^i, \tilde{\bar{Z}}_{n'}^{\bar{j}}] = n\delta_{n+n'}\eta^{i\bar{j}}\,. \tag{A.3}$$

The oscillator vacuum is defined by $Z_n|0\rangle = 0, \tilde{Z}_n|0\rangle = 0$ for $n > 0$.

A.2 Internal bosons — 11 sector

For open strings, the worldsheet coordinate $\sigma \in [0, \pi]$. In the 11 sector, the open strings satisfy the boundary conditions

$$\operatorname{Im} Z^i|_{\sigma=0,\pi} = 0\,; \qquad \operatorname{Re}(\partial_\sigma Z^i)|_{\sigma=0,\pi} = 0\,. \tag{A.4}$$

In terms of oscillator modes, this implies that $Z_n^i = -\tilde{\bar{Z}}_n^{\bar{i}}$, and $\operatorname{Im} p^i = 0$. The hamiltonian L_0 for these modes is:

$$L_0^{\mathrm{bos}} = \frac{1}{4}\sum_i(\operatorname{Re} p^i)^2 + \sum_{i,\ n^i>0}\left(\bar{Z}_{-n^i}^{\bar{i}} Z_{n^i}^i + Z_{-n^i}^i \bar{Z}_{n^i}^{\bar{i}}\right). \tag{A.5}$$

A.3 Internal bosons — 22 sector

In the 22 sector, the open strings satisfy the boundary conditions

$$\operatorname{Im}\left(e^{i\theta_i}Z^i|_{\sigma=0,\pi}\right) = 0\,; \qquad \operatorname{Re}\left(e^{i\theta_i}\partial_\sigma Z^i\right)|_{\sigma=0,\pi} = 0\,. \tag{A.6}$$

In terms of oscillator modes, this implies that $Z_n^i = -e^{-2i\theta_i}\tilde{\bar{Z}}_n^{\bar{i}}$, and $\operatorname{Im} p^i = 0$. The hamiltonian L_0 for these modes is:

$$L_0^{\mathrm{bos}} = \frac{1}{4}\sum_i^3(\operatorname{Re} p^i)^2 + \sum_{i;\ n^i>0}\left(\bar{Z}_{-n^i}^{\bar{i}} Z_{n^i}^i + Z_{-n^i}^i \bar{Z}_{n^i}^{\bar{i}}\right). \tag{A.7}$$

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A.4 Internal bosons — 12 and 21 sectors

In the 12 sector, the strings end on brane 1 at $\sigma = 0$ and on brane 2 at $\sigma = \pi$. Therefore,

$$\begin{aligned} \operatorname{Im} Z^i|_{\sigma=0} &= 0\,; \qquad \operatorname{Re}(\partial_\sigma Z^i)|_{\sigma=0} = 0 \operatorname{Im}\left(e^{i\theta_i} Z^i|_{\sigma=\pi}\right) \\ &= 0\,; \qquad \operatorname{Re}\left(e^{i\theta_i}\partial_\sigma Z^i\right)|_{\sigma=\pi} = 0\,. \end{aligned} \tag{A.8}$$

In order to satisfy these boundary conditions, we must change the moding of the oscillators. The mode expansion in this sector is:

$$\begin{aligned} Z^i &= z_0 + \frac{i}{\sqrt{4\pi}} \sum_{n=-\infty}^{\infty} \left(\frac{Z^i_{n+\alpha_i}}{n+\alpha_i} e^{-i(n+\alpha_i)(\tau-\sigma)} + \frac{\tilde{Z}^i_{n-\alpha_i}}{n-\alpha_i} e^{-i(n-\alpha_i)(\tau+\sigma)} \right) \\ \bar{Z}^{\bar{i}} &= z_0 + \frac{i}{\sqrt{4\pi}} \sum_{n=-\infty}^{\infty} \left(\frac{\bar{Z}^{\bar{i}}_{n+\alpha_i}}{n+\alpha_i} e^{-i(n+\alpha_i)(\tau-\sigma)} + \frac{\tilde{\bar{Z}}^{\bar{i}}_{n-\alpha_i}}{n-\alpha_i} e^{-i(n-\alpha_i)(\tau+\sigma)} \right) \end{aligned} \tag{A.9}$$

where $\alpha_i = \theta_i/\pi$. Canonical commutation relations imply the following commutators for the modes:

$$[Z^i_{n+\alpha_i}, \bar{Z}^{\bar{j}}_{n'+\alpha_{\bar{j}}}] = (n+\alpha_i)\delta_{n+n'}\eta^{i\bar{j}}\,. \tag{A.10}$$

The boundary conditions (A.8) can then be written in terms of the modes as:

$$Z^i_{n+\alpha_i} = -\tilde{\bar{Z}}^{\bar{i}}_{n+\alpha_i}\,, \tag{A.11}$$

and the zero modes are forced to vanish, as one might expect: motion away from this intersection forces the string to stretch out, and is not a zero mode. The hamiltonian for modes of Z^i in this sector is:

$$\begin{aligned} L_0^{12,\alpha>0} &= \sum_{n\geq 0} \bar{Z}^{\bar{i}}_{-n+\alpha} Z^i_{n+\alpha} + \sum_{n>0} Z^i_{-n+\alpha_i} \bar{Z}^{\bar{i}}_{n+\alpha} \qquad \text{if } \alpha > 0 \\ L_0^{12,\alpha<0} &= \sum_{n>0} \bar{Z}^{\bar{i}}_{-n+\alpha} Z^i_{n+\alpha} + \sum_{n\geq 0} Z^i_{-n+\alpha_i} \bar{Z}^{\bar{i}}_{n+\alpha} \qquad \text{if } \alpha < 0\,. \end{aligned} \tag{A.12}$$

Using the canonical commutation relations, the spectrum of L_0 in this sector is:

$$\begin{aligned} E_{N,\bar{N}} &= \sum_{n\geq 0}(n+\alpha)N_n + \sum_{\bar{n}>0}(\bar{n}-\alpha)\bar{N}_{\bar{n}} \quad \text{if } \alpha > 0 \\ E_{N,\bar{N}} &= \sum_{n>0}(n+\alpha)N_n + \sum_{\bar{n}\geq 0}\bar{n}-\alpha)\bar{N}_{\bar{n}} \quad \text{if } \alpha < 0 \end{aligned} \tag{A.13}$$

where $N_n, \bar{N}_{\bar{n}} \in \mathbb{Z}$.

A.5 Internal fermions — closed string sector

The fermions $\psi^i_\pm, \bar{\psi}^{\bar{i}}_\pm$ are the worldsheet superpartners of $Z^i, \bar{Z}^{\bar{i}}$. The mode expansions in the closed string sector are:

$$\begin{aligned} \psi^i_+ &= \sum_n \psi^i_{+,n} e^{-in(\tau+\sigma)}\,; \qquad \bar{\psi}^{\bar{i}}_+ = \sum_n \bar{\psi}^{\bar{i}}_{+,n} e^{-in(\tau+\sigma)} \\ \psi^i_- &= \sum_n \psi^i_{-,n} e^{-in(\tau-\sigma)}\,; \qquad \bar{\psi}^{\bar{i}}_- = \sum_n \bar{\psi}^{\bar{i}}_{-,n} e^{-in(\tau-\sigma)}\,. \end{aligned} \tag{A.14}$$

The $\mathrm{U}(1)_R$ currents are $J_{R,\pm} = \eta_{i\bar{\jmath}}\psi^i_\pm\bar{\psi}^{\bar{\jmath}}_\pm$. Fermions $\psi^i_\pm$ have charge $+1$ under this $\mathrm{U}(1)_R$, while the fermions $\bar{\psi}^{\bar{\jmath}}_\pm$ have charge -1. For NS fermions, $n \in \mathbb{Z}+\frac{1}{2}$, while for R fermions, $n \in \mathbb{Z}$. The canonical anticommutation relations for $\psi(\sigma)$ imply the following anticommutation relations for the modes:

$$\{\psi^i_{\pm,n}, \bar{\psi}^{\bar{\jmath}}_{\pm,n'}\} = \delta_{n+n'}\eta^{i\bar{\jmath}}\,. \tag{A.15}$$

The vacuum is defined by $\psi^i_{\pm,n}|0\rangle = 0$ for $n > 0$. For R fermions, the $n = 0$ modes form a Clifford algebra:

$$\{\psi^i_{\pm,0}, \bar{\psi}^{\bar{\jmath}}_{\pm,0}\} = \eta^{i\bar{\jmath}}\,. \tag{A.16}$$

For a given complex direction i, this algebra has a two-dimensional representation. For example, for $\psi^1_{0,+}, \bar{\psi}^{\bar{1}}_{0,+}$, we can write

$$\begin{aligned}
\bar{\psi}^{\bar{1}}_{0,+}|\downarrow_{1,+}\rangle &= 0\\
\psi^1_{0,+}|\downarrow_{1,+}\rangle &= |\uparrow_{1,+}\rangle\\
\psi^1_{0,+}|\uparrow_{1,+}\rangle &= 0
\end{aligned} \tag{A.17}$$

CPT invariance requires that $|\uparrow_{i,\pm}\rangle$ and $|\downarrow_{i,\pm}\rangle$ have R-charges $+\frac{1}{2}$ and $-\frac{1}{2}$, respectively.

A.6 Internal fermions — 11 sector

The boundary conditions for the fermions are related via supersymmetry to those of the bosons. In the 11 sector, these conditions are:

$$\psi^i_+ = \bar{\psi}^{\bar{i}}_-\,. \tag{A.18}$$

at both $\sigma = 0, \pi$. In terms of modes, this implies:

$$\psi^i_{+,n} = \bar{\psi}^{\bar{i}}_{-,n}\,. \tag{A.19}$$

The worldsheet hamiltonian for these modes is:

$$L_0 = \sum_{n>0} n\left(\psi^{\bar{i}}_{+,-n}\psi^i_{+,n} + \psi^i_{+,-n}\psi^{\bar{i}}_{+,n}\right). \tag{A.20}$$

A.7 Internal fermions — 22 sector

In the 22 sector, the boundary conditions are:

$$\psi^i_+ = e^{-2i\theta_i}\bar{\psi}^{\bar{i}}_-\,. \tag{A.21}$$

at both $\sigma = 0, \pi$. In terms of modes, this implies:

$$\psi^i_{+,n} = e^{-2i\theta_i}\bar{\psi}^{\bar{i}}_{-,n}\,. \tag{A.22}$$

The worldsheet hamiltonian is the same as (A.20).

A.8 Internal fermions — 12 and 21 sectors

Here the boundary conditions are

$$\begin{aligned}\psi^i_+(\sigma=0) &= \bar{\psi}^{\bar{i}}_-(\sigma=0)\\ \psi^i_+(\sigma=\pi) &= e^{-2i\theta_i}\bar{\psi}^{\bar{i}}_-(\sigma=\pi)\end{aligned}\tag{A.23}$$

To solve this, we must shift the moding:

$$\begin{aligned}\psi^i_\pm &= \sum_n \psi^i_{\pm,n\pm\alpha_i} e^{-i(n\pm\alpha_i)(\tau\pm\sigma)}\\ \bar{\psi}^{\bar{i}}_\pm &= \sum_n \bar{\psi}^{\bar{i}}_{\pm,n\pm\alpha_i} e^{-i(n\pm\alpha_i)(\tau\pm\sigma)}\end{aligned}\tag{A.24}$$

where $n \in \mathbb{Z}+\frac{1}{2}$ for NS fermions, and $n \in \mathbb{Z}$ for R fermions. The anticommutation relations for the modes are:

$$\{\psi^i_{\pm,n\pm\alpha_i}, \bar{\psi}^{\bar{j}}_{\pm,n'\pm\alpha_i}\} = \eta^{i\bar{j}}\delta_{n+n'}\,,\tag{A.25}$$

and the boundary conditions are:

$$\begin{aligned}\psi^i_{+,n+\alpha_i} &= \bar{\psi}^{\bar{i}}_{-,n-\alpha_i}\\ \bar{\psi}^{\bar{i}}_{+,n+\alpha_i} &= \psi^i_{-,n-\alpha_i}\,.\end{aligned}\tag{A.26}$$

The hamiltonian in this sector is:

$$L_0 = \sum_{n>-\alpha_i}(n+\alpha_i)\bar{\psi}^{\bar{i}}_{+,-n+\alpha_i}\psi^i_{+,n+\alpha_i} + \sum_{n>\alpha_i}(n-\alpha_i)\psi^i_{-,-n+\alpha_i}\bar{\psi}^{\bar{i}}_{+,n+\alpha_i}\tag{A.27}$$

with the spectrum

$$E_{N,\bar{N}} = \sum_{n>-\alpha_i}(n+\alpha_i)N_n + \sum_{n>\alpha_i}(n-\alpha_i)\bar{N}_{\bar{n}}\,.\tag{A.28}$$

Here $N, \bar{N} = 0$ or 1.

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