PROCEEDINGS OF SCIENCE

Light strings at D-brane intersections

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This proceeding is based on [1] where we study the spectrum of open strings localized at the intersections of D6-branes. For a regime in parameter space we find that the masses of the lightest states scale as $M_\theta^2 \approx \theta M_s^2$ and can thus be parametrically smaller than the string scale. Relying on previous analyses, we compute scattering amplitudes of massless 'twisted' open strings and study their factorization, confirming the presence of the light massive states as sub-dominant poles in one of the channels. Along the analysis we provide a dictionary between states and their corresponding vertex operators.

Proceedings of the Corfu Summer Institute 2011 School and Workshops on Elementary Particle Physics and Gravity
September 4-18, 2011
Corfu, Greece

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PoS(CORFU2011)030

1. Introduction

D-brane compactifications provide a promising framework for model building [2–5]. One of the peculiar features is the possibility of accommodating large extra dimensions giving rise to a significantly lower string scale, even of a few *TeV* [6–8]. Scenarios of these kinds may circumvent the hierarchy problem, and also allow for stringy signatures that can be observed at LHC [9–22].

Recently, in a series of papers [23–26] the authors study tree-level string scattering amplitudes containing at most two chiral fermions. They show that these amplitudes exhibit a universal behaviour independently of the specifics of the compactification, which gives their results a predictive power. The observed poles correspond to the exchanges of Regge excitations of the standard model gauge bosons, whose masses scale with the string mass M_s.

On the other hand there exists a tower of stringy excitations of the chiral fermions and their superpartners localized at the intersections of two stacks of D-branes. Their masses depend on the string mass M_s and the intersection angle θ and thus can be significantly lighter than the Regge excitations of the gauge bosons, given that the intersection angle θ is small. In this note we will analyse in detail the spectrum arising at the intersection of two D-branes. We will see that for a regime in parameter space some of those states can be very light and furthermore discuss how such states can be potentially observed.

2. Quantization at Angles

Let us start by solving the equations of motion for an open string stretched between two D-brane stacks intersecting at an angle $\pi\theta$ in the (X,Y) plane. The bosonic coordinates have to fulfil the boundary conditions [27–30]

$$\begin{aligned} \partial_\sigma X(\tau,0) &= 0 = Y(\tau,0) \\ \partial_\sigma X(\tau,\pi) + \tan(\pi\theta)\,\partial_\sigma Y(\tau,\pi) &= 0 \\ Y(\tau,\pi) - \tan(\pi\theta)\,X(\tau,\pi) &= 0\,. \end{aligned} \tag{2.1}$$

It proves convenient to introduce complex coordinates $Z^I = X^I + \mathrm{i}Y^I$ with $I = 1,2,3$ for the internal (compactified) directions. Given these boundary conditions for each X and Y, one can deduce the mode expansions for each ∂Z and $\partial\overline{Z}$ that read (after applying the doubling trick)

$$\partial Z(z) = \sum_n \alpha_{n-\theta}\, z^{-n+\theta-1} \qquad\qquad \partial\overline{Z}(z) = \sum_n \alpha_{n+\theta}\, z^{-n-\theta-1}\,. \tag{2.2}$$

Upon quantization the only non-vanishing commutators are

$$[\alpha_{n\pm\theta}, \alpha_{m\mp\theta}] = (m\pm\theta)\,\delta_{n+m}\,.$$

World-sheet supersymmetry $\delta X = \overline{\varepsilon}\psi$ leads to the same modding for the complexified world-sheet fermions. One obtains (again after using the doubling trick)

$$\Psi(z) = \sum_{r\in\mathbb{Z}+\nu} \psi_{r-\theta}\, z^{-r-\frac{1}{2}+\theta} \qquad \overline{\Psi}(z) = \sum_{r\in\mathbb{Z}+\nu} \psi_{r+\theta}\, \overline{z}^{-r-\frac{1}{2}-\theta}\,, \tag{2.3}$$

where ν is $\frac{1}{2}$ and 0 for the NS-sector and R-sector, respectively. Upon quantization the only non-vanishing anti-commutators are

$$\{\psi_{m-\theta}, \psi_{n+\theta}\} = \delta_{m,n}\,. \tag{2.4}$$

2.1 NS-sector

Let us start with the NS sector that describes space-time bosons restricting for the moment our attention onto just one complex dimension. The definition of the ground-state crucially depends on whether the intersection angles are positive or negative. For a positive intersection angle the ground-state $|\,\theta_I\,\rangle_{NS}$ is given by

$$\begin{aligned} &\alpha_{m-\theta}|\,\theta\,\rangle_{NS}=0 \qquad m\geq 1 \qquad\qquad \psi_{r-\theta}|\,\theta\,\rangle_{NS}=0 \qquad r\geq\frac{1}{2}\\ &\alpha_{m+\theta}|\,\theta\,\rangle_{NS}=0 \qquad m\geq 0 \qquad\qquad \psi_{r+\theta}|\,\theta\,\rangle_{NS}=0 \qquad r\geq\frac{1}{2}\end{aligned} \tag{2.5}$$

whereas for a negative intersection angle it is defined as

$$\begin{aligned} &\alpha_{m-\theta}|\,\theta\,\rangle_{NS}=0 \qquad m\geq 0 \qquad\qquad \psi_{r-\theta}|\,\theta\,\rangle_{NS}=0 \qquad r\geq\frac{1}{2}\\ &\alpha_{m+\theta}|\,\theta\,\rangle_{NS}=0 \qquad m\geq 1 \qquad\qquad \psi_{r+\theta}|\,\theta\,\rangle_{NS}=0 \qquad r\geq\frac{1}{2}\,.\end{aligned} \tag{2.6}$$

Due to the non-trivial intersection angles the vertex operators describing the states under consideration involve bosonic and fermionic twist fields accounting for the boundary conditions (2.1). In order to properly identify these twist fields we determine the action of the conformal fields Ψ, $\overline{\Psi}$, ∂Z and $\partial\overline{Z}$ on the ground-state $|\,\theta\,\rangle_{NS}$ and excitations (fermionic and bosonic ones) thereof and obtain the following dictionary [1].

Positive angles		Negative angles	
state	vertex operator	state	vertex operator
$\|\,\theta\,\rangle_{NS}$	$e^{i\theta H(z)}\sigma^+_\theta(z)$	$\|\,\theta\,\rangle_{NS}$	$e^{i\theta H(z)}\sigma^-_\theta(z)$
$\alpha_{-\theta}\|\,\theta\,\rangle_{NS}$	$e^{i\theta H(z)}\tau^+_\theta(z)$	$\alpha_\theta\|\,\theta\,\rangle$	$e^{i\theta H(z)}\widetilde{\tau}^-_\theta(z)$
$(\alpha_{-\theta})^2\|\,\theta\,\rangle_{NS}$	$e^{i\theta H(z)}\omega^+_\theta(z)$	$(\alpha_\theta)^2\|\,\theta\,\rangle_{NS}$	$e^{i\theta H(z)}\widetilde{\omega}^-_\theta(z)$
$\psi_{-\frac{1}{2}+\theta}\|\,\theta\,\rangle_{NS}$	$e^{i(\theta-1)H(z)}\sigma^+_\theta(z)$	$\psi_{-\frac{1}{2}-\theta}\|\,\theta\,\rangle_{NS}$	$e^{i(\theta+1)H(z)}\sigma^-_\theta(z)$
$\alpha_{-\theta}\,\psi_{-\frac{1}{2}+\theta}\|\,\theta\,\rangle_{NS}$	$e^{i(\theta-1)H(z)}\tau^+_\theta(z)$	$\alpha_\theta\,\psi_{-\frac{1}{2}-\theta}\|\,\theta\,\rangle_{NS}$	$e^{i(\theta+1)H(z)}\widetilde{\tau}^-_\theta(z)$
$(\alpha_{-\theta})^2\,\psi_{-\frac{1}{2}+\theta}\|\,\theta\,\rangle_{NS}$	$e^{i(\theta-1)H(z)}\omega^+_\theta(z)$	$(\alpha_\theta)^2\,\psi_{-\frac{1}{2}-\theta}\|\,\theta\,\rangle_{NS}$	$e^{i(\theta+1)H(z)}\widetilde{\omega}^-_\theta(z)$
$\alpha_{-1+\theta}\|\,\theta\,\rangle_{NS}$	$e^{i\theta H(z)}\widetilde{\tau}^+_\theta(z)$	$\alpha_{-1-\theta}\|\,\theta\,\rangle_{NS}$	$e^{i\theta H(z)}\tau^-_\theta(z)$
$\alpha_{-1+\theta}\psi_{-\frac{1}{2}+\theta}\|\,\theta\,\rangle_{NS}$	$e^{i(\theta-1)H(z)}\widetilde{\tau}^+_\theta(z)$	$\alpha_{-1-\theta}\psi_{-\frac{1}{2}-\theta}\|\,\theta\,\rangle_{NS}$	$e^{i(\theta+1)H(z)}\tau^-_\theta(z)$

Table 1: Excitations and their corresponding vertex operator part for the NS-sector.

2.2 R-sector

In the R-sector the bosonic part behaves exactly the same as in the NS-sector. Thus it is sufficient to study the fermionic part. The mode expansion of Ψ and $\overline{\Psi}$ are similar to the expansions in the NS sector however the sum is over integers and not half-integers (see eq. (2.3)). Again the definition of the ground state crucially depends on whether the intersection angle is positive or negative. For positive intersection angle one has

$$\begin{aligned} &\alpha_{m-\theta}|\,\theta\,\rangle_{R}=0 \qquad m\geq 1 \qquad\qquad \psi_{r-\theta}|\,\theta\,\rangle_{R}=0 \qquad r\geq 1\\ &\alpha_{m+\theta}|\,\theta\,\rangle_{R}=0 \qquad m\geq 0 \qquad\qquad \psi_{r+\theta}|\,\theta\,\rangle_{R}=0 \qquad r\geq 0\end{aligned} \tag{2.7}$$

whereas for a negative intersection angle one defines

$$\begin{array}{llll} \alpha_{m-\theta}|\,\theta\,\rangle_R = 0 & m \geq 0 & \qquad \psi_{r-\theta}|\,\theta\,\rangle_R = 0 & r \geq 0 \\ \alpha_{m+\theta}|\,\theta\,\rangle_R = 0 & m \geq 1 & \qquad \psi_{r+\theta}|\,\theta\,\rangle_R = 0 & r \geq 1\;. \end{array} \tag{2.8}$$

Due to the fact that the mode expansion of Ψ and $\overline{\Psi}$ in the R-sector is over integers rather than half-integers the fermionic twist operators will take a different form from the ones in the NS-sector. Applying the same procedure as in the NS-sector we obtain.

Positive angles		Negative angles	
state	vertex operator	state	vertex operator
$\|\,\theta\,\rangle_R$	$e^{i\left(\theta-\frac{1}{2}\right)H(z)}\sigma^+_\theta(z)$	$\|\,\theta\,\rangle_R$	$e^{i\left(\frac{1}{2}-\theta\right)H(z)}\sigma^-_\theta(z)$
$\alpha_{-\theta}\|\,\theta\,\rangle_R$	$e^{i\left(\theta-\frac{1}{2}\right)H(z)}\tau^+_\theta(z)$	$\alpha_\theta\|\,\theta\,\rangle_R$	$e^{i\left(\frac{1}{2}-\theta\right)H(z)}\widetilde{\tau}^-_\theta(z)$
$\psi_{-\theta}\|\,\theta\,\rangle_R$	$e^{i\left(\theta+\frac{1}{2}\right)H(z)}\sigma^+_\theta(z)$	$\psi_\theta\|\,\theta\,\rangle_R$	$e^{i\left(\theta-\frac{1}{2}\right)H(z)}\sigma^-_\theta(z)$
$\alpha_{-\theta}\,\psi_{-\theta}\|\,\theta\,\rangle_R$	$e^{i\left(\theta+\frac{1}{2}\right)H(z)}\tau^+_\theta(z)$	$\alpha_\theta\,\psi_\theta\|\,\theta\,\rangle_R$	$e^{i\left(\theta-\frac{1}{2}\right)H(z)}\widetilde{\tau}^-_\theta(z)$
$\psi_{-1+\theta}\|\,\theta\,\rangle_R$	$e^{i\left(\theta-\frac{3}{2}\right)H(z)}\sigma^+_\theta(z)$	$\psi_{-1-\theta}\|\,\theta\,\rangle_R$	$e^{i\left(\theta+\frac{3}{2}\right)H(z)}\sigma^-_\theta(z)$
$\alpha_{-\theta}\,\psi_{-1+\theta}\|\,\theta\,\rangle_R$	$e^{i\left(\theta-\frac{3}{2}\right)H(z)}\tau^+_\theta(z)$	$\alpha_\theta\,\psi_{-1-\theta}\|\,\theta\,\rangle_R$	$e^{i(\left(\theta+\frac{3}{2}\right)H(z)}\widetilde{\tau}^-_\theta(z)$

Table 2: Excitations and their corresponding vertex operator part for the R-sector.

3. D-brane setup

Let us briefly recall the main features of intersecting brane worlds [2–5]. The gauge groups arise from stacks of D6-branes that fill out four- dimensional space-time and wrap three-cycles in the internal Calabi-Yau threefold. Chiral matter appears at the intersection in the internal space of different cycles wrapped by the D6-brane stacks. The multiplicity of chiral matter between two stacks of D6-branes is given by the topological intersection number of the respective three-cycles.

Many features of a D-brane compactifications, such as chiral matter, gauge symmetry or Yukawa couplings do not crucially depend on the details of the compactification, but rather only on the local structure of the D-brane configurations. Thus it is often times sufficient to investigate a local D-brane setup, described by some quiver theory, and to postpone the embedding into a global setting. This approach is called bottom-up approach and has been initiated in [31,32][1].

In the following analysis we have in mind such a local D-brane configuration. However, instead of looking at the whole local configuration we further zoom in and just focus on a subset of the D-brane stacks and investigate the various states localized at the intersection of two stacks. Let us further specify the setup. We have three stacks of D6-branes wrapping three-cycles on the factorizable six-torus $T^6 = T^2 \times T^2 \times T^2$ [29,37,38]. They intersect each other non-trivially and

[1]For a systematic search of realistic MSSM D-brane quivers, see [33, 34]. For an exhaustive search of global embeddings of such quivers, see [35,36].

give rise to the following intersection angles

$$\begin{array}{lll} \theta^1_{ab} > 0 & \theta^2_{ab} > 0 & \theta^3_{ab} < 0 \\ \theta^1_{bc} > 0 & \theta^2_{bc} > 0 & \theta^3_{bc} < 0 \\ \theta^1_{ca} < 0 & \theta^2_{ca} < 0 & \theta^3_{ca} < 0\,. \end{array} \tag{3.1}$$

At each intersection massless chiral fermions appear and, in case of a preserved supersymmetry,

$$\theta^1_{ab} + \theta^2_{ab} + \theta^3_{ab} = 0 \qquad \theta^1_{bc} + \theta^2_{bc} + \theta^3_{bc} \quad = 0 \qquad \theta^1_{ca} + \theta^2_{ca} + \theta^3_{ca} = -2 \tag{3.2}$$

even massless scalars. We do not always have to enforce the supersymmetry constraints, since the analysis applies independently of whether supersymmetry is preserved or not. Moreover, apart from the massless matter at each intersection there are also massive states whose mass scales with the intersection angle. In scenarios with a low string tension and small intersection angles such states can be fairly light and potentially observed at LHC or future experiments. Let us take a look at each intersection separately and assume without loss of generality that the intersection angle θ^1_{ca} is small.

We start with the massless chiral fermions appearing at the intersection ab and bc, given by the Ramond vacuum. Applying the procedure laid out above to the choice of intersection angles (3.1) one obtains for the vertex operator of $|\,\theta_{1,2,3}\,\rangle^{ab}_R$ [2]

$$|\,\theta_{1,2,3}\,\rangle^{ab}_R : \qquad V^{(-1/2)}_{\psi} = \Lambda_{ab}\psi^{\alpha}\, e^{-\varphi/2} S_{\alpha} \prod_{I=1}^{2} \sigma^{+}_{\theta^I_{ab}}\, e^{i\left(\theta^I_{ab}-\frac{1}{2}\right)H_I}\, \sigma^{-}_{-\theta^3_{ab}}\, e^{i\left(\theta^3_{ab}+\frac{1}{2}\right)H_3}\, e^{ikX}\,. \tag{3.3}$$

Its right-handed counterpart is given by

$$|\,\theta_{1,2,3}\,\rangle^{ba}_R : \qquad V^{(-1/2)}_{\overline{\psi}} = \Lambda_{ba}\overline{\psi}_{\dot\alpha}\, e^{-\varphi/2} S^{\dot\alpha} \prod_{I=1}^{2} \sigma^{-}_{\theta^I_{ab}}\, e^{i\left(-\theta^I_{ab}+\frac{1}{2}\right)H_I}\, \sigma^{+}_{-\theta^3_{ab}}\, e^{i\left(-\theta^3_{ab}-\frac{1}{2}\right)H_3}\, e^{ikX}\,. \tag{3.4}$$

Similarly we get for the bc sector

$$|\,\theta_{1,2,3}\,\rangle^{bc}_R : \qquad V^{(-1/2)}_{\chi} = \Lambda_{bc}\chi^{\alpha}\, e^{-\varphi/2} S_{\alpha} \prod_{I=1}^{2} \sigma^{+}_{\theta^I_{bc}}\, e^{i\left(\theta^I_{bc}-\frac{1}{2}\right)H_I}\, \sigma^{-}_{-\theta^3_{bc}}\, e^{i\left(\theta^3_{bc}+\frac{1}{2}\right)H_3}\, e^{ikX} \tag{3.5}$$

and

$$|\,\theta_{1,2,3}\,\rangle^{cb}_R : \qquad V^{(-1/2)}_{\overline{\chi}} = \Lambda_{cb}\overline{\chi}_{\dot\alpha}\, e^{-\varphi/2} S^{\dot\alpha} \prod_{I=1}^{2} \sigma^{-}_{\theta^I_{bc}}\, e^{i\left(-\theta^I_{bc}+\frac{1}{2}\right)H_I}\, \sigma^{+}_{-\theta^3_{bc}}\, e^{i\left(-\theta^3_{bc}-\frac{1}{2}\right)H_3}\, e^{ikX}\,. \tag{3.6}$$

Let us now turn to the ca sector, where one can observe light stringy states. The massless chiral fermion is given again by the Ramond vacuum $|\,\theta_{1,2,3}\,\rangle^{ca}_R$

$$|\,\theta_{1,2,3}\,\rangle^{ca}_R : \qquad V^{(-1/2)}_{\psi} = \Lambda_{ca}\psi^{\alpha}\, e^{-\varphi/2} S_{\alpha} \prod_{I=1}^{3} \sigma^{-}_{\theta^I_{ca}}\, e^{i\left(\theta^I_{ab}+\frac{1}{2}\right)H_I}\, e^{ikX}\,. \tag{3.7}$$

[2] For a discussion of vertex operators for massless states at arbitrary intersection angles, see [39,40]. For an analysis of instantonic modes at the intersection of D-instanton and D-brane at arbitrary angles, see [41]. Vertex operators of massive states in heterotic compactifications are discussed in [42]

In the NS-sector we have the massless (if SUSY is preserved) scalar $\prod_{I=1}^{3}\psi^{I}_{-\frac{1}{2}-\theta^{I}_{ca}}\,|\,\theta_{1,2,3}\,\rangle^{ca}_{NS}$ whose vertex operator is given by

$$V^{(-1)}_{\phi}=\Lambda_{ca}\,\phi\;e^{-\varphi}\prod_{I=1}^{3}\sigma^{-}_{\theta^{I}_{ca}}\;e^{i(1+\theta^{I}_{ca})H_I}\;e^{ikX} \tag{3.8}$$

The first bosonic excitation is $\alpha^{1}_{-\theta^{1}_{ca}}\prod_{I=1}^{3}\psi^{I}_{-\frac{1}{2}-\theta^{I}_{ca}}\,|\,\theta_{1,2,3}\,\rangle^{ca}_{NS}$ and has mass $M^2=-\theta^{1}_{ca}M^2_s$. Its vertex operator is

$$V^{(-1)}_{\tilde\phi}=\Lambda_{ca}\,\tilde\phi\;e^{-\varphi}\;\widetilde{\tau}^{-}_{\theta^{1}_{ca}}\;e^{i(1+\theta^{1}_{ca})H_1}\prod_{I=2}^{3}\sigma^{-}_{\theta^{I}_{ca}}\;e^{i(1+\theta^{I}_{ca})H_I}\;e^{ikX}\;. \tag{3.9}$$

The second bosonic excitation is $\left(\alpha^{1}_{-\theta^{1}_{ca}}\right)^{2}\prod_{I=1}^{3}\psi^{I}_{-\frac{1}{2}-\theta^{I}_{ca}}\,|\,\theta_{1,2,3}\,\rangle^{ca}_{NS}$, that has mass $M^2=-2\theta^{1}_{ca}M^2_s$ and whose vertex operator takes the form

$$V^{(-1)}_{\widehat\phi}=\Lambda_{ca}\,\widehat\phi\;e^{-\varphi}\;\widetilde{\omega}^{-}_{\theta^{1}_{ca}}\;e^{i(1+\theta^{1}_{ca})H_1}\prod_{I=2}^{3}\sigma^{-}_{\theta^{I}_{ca}}\;e^{i(1+\theta^{I}_{ca})H_I}\;e^{ikX}\;. \tag{3.10}$$

It is easy to check that the conformal dimensions of these vertex operators indeed account for states with mass $M^2=-\theta^{1}_{ca}M^2_s$ and $M^2=-2\theta^{1}_{ca}M^2_s$, respectively. Given the fact that the intersection angle θ^{1}_{ca} is small and assuming a low string scale these states may be observable. In the following we study the four point amplitude $\langle\overline{\psi}\,\psi\,\chi\,\overline{\chi}\rangle$ which in the t-channel exposes such light states.

3.1 The amplitude

Here we compute the scattering amplitude of four chiral fermions $\langle\overline{\psi}\,\psi\,\chi\,\overline{\chi}\rangle$ where ψ and χ are the chiral massless fermions localized at the intersection *ab*, and *bc*, respectively. The fields $\overline{\psi}$ and $\overline{\chi}$ are their corresponding anti-particle. Before giving the details of the computation let us discuss briefly the naive expectations concerning various limits of this amplitude.

In the *s*-channel, displayed in figure 1a, one expects the exchange of a gauge boson living on the D-brane stack *b*. Indeed the dominant pole indicates a gauge boson exchange that allows one to normalize the four-point amplitude. Higher poles correspond to exchanges of stringy excitations whose masses scale as M_s. Such states can already be observed in the scattering amplitude of four gauge bosons and also in scattering of two fermions onto two gauge bosons. For a sufficiently small string tension, in the *TeV* range, one may observe signals of these states at LHC [23, 25].

On the other hand in the *t*-channel, displayed in figure 1b, the dominant pole indicates the exchange of a scalar which is massless if supersymmetry is preserved. Furthermore one expects additional poles corresponding to exchanges of massive stringy states. In contrast to the *s*-channel exchange particles the masses of those states do not only scale with M_s but also with the intersection angle θ_{ac}. Thus they could be significantly lighter for small intersection angle θ_{ac} and signals of such states are expected to be observed even before observations of the massive untwisted stringy states.

Given the vertex operators (3.3) to (3.6) we are now able to compute the amplitude

$$\mathscr{A}=\langle\overline{\psi}(0)\;\psi(x)\;\chi(1)\;\overline{\chi}(\infty)\rangle \tag{3.11}$$

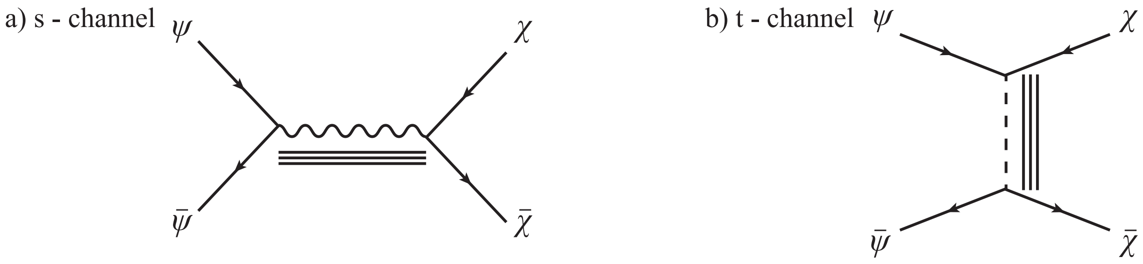

Figure 1: The s-channel: the curly line denotes the gauge boson. The t-channel: the dashed line denotes the massless scalar. The solid lines denote massive stringy states.

and obtain

$$\mathscr{A} \sim i g_s \,\mathrm{Tr}\left(\Lambda_{ba}\,\Lambda_{ab}\,\Lambda_{bc}\,\Lambda_{cb}\right) \overline{\psi}\cdot\overline{\chi}\,\psi\cdot\chi\,(2\pi)^4\delta^{(4)}\left(\sum_i^4 k_i\right) \tag{3.12}$$
$$\times \int_0^1 dx \frac{x^{-1+k_1\cdot k_2}\,(1-x)^{-\frac{3}{2}+k_2\cdot k_3}\, e^{-S_{cl}(\theta^1_{ab},1-\theta^1_{bc})}\, e^{-S_{cl}(\theta^2_{ab},1-\theta^2_{bc})}\, e^{-S_{cl}(1+\theta^3_{ab},-\theta^3_{bc})}}{\left[I(\theta^1_{ab},1-\theta^1_{bc},x)\,I(\theta^2_{ab},1-\theta^2_{bc},x)\,I(1+\theta^3_{ab},-\theta^3_{bc},x)\right]^{\frac{1}{2}}}\,.$$

where we used the standard correlators. For details we refer the interested reader to [1].

s-channel – normalization of the amplitude

The s-channel allows us to normalize the amplitude. In order to properly take the limit $x \to 0$ we Poissón resum the classical contribution, obtaining

$$\mathscr{A} \sim i g_s \mathrm{Tr}\left(\Lambda_{ba}\,\Lambda_{ab}\,\Lambda_{bc}\,\Lambda_{cb}\right) \overline{\psi}\cdot\overline{\chi}\,\psi\cdot\chi\,\frac{(2\pi)^4\delta^{(4)}\left(\sum_i^4 k_i\right)}{L_{b^1}\,L_{b^2}\,L_{b^3}} \tag{3.13}$$
$$\times \int_0^1 dx \frac{x^{-1+k_1\cdot k_2}\,(1-x)^{-\frac{3}{2}+k_2\cdot k_3}\, e^{-\widetilde{S}_{cl}(\theta^1_{ab},1-\theta^1_{bc})}\, e^{-\widetilde{S}_{cl}(\theta^2_{ab},1-\theta^2_{bc})}\, e^{-\widetilde{S}_{cl}(1+\theta^3_{ab},-\theta^3_{bc})}}{\sqrt{{}_2F_1[\theta^1_{ab},\theta^1_{bc},1;x]\,{}_2F_1[\theta^2_{ab},\theta^2_{bc},1;x]\,{}_2F_1[1+\theta^3_{ab},1+\theta^3_{bc},1;x]}}\,,$$

where $e^{-\widetilde{S}_{cl}}$ in the Hamiltonian form is given by [30, 43–47]

$$e^{-\widetilde{S}_{cl}(\theta,\nu)} = \prod_{i=1}^{3}\sum_{p_i,q_i} \exp\left[-\pi\frac{t(\theta,\nu,x)}{\sin(\pi\theta)}\,\frac{\alpha'}{L^2_{b_i}}\,p_i^2 - \pi\frac{t(\theta,\nu,x)}{\sin(\pi\theta)}\,\frac{R^2_{x_i}\,R^2_{y_i}}{\alpha'\,L^2_{b^i}}\,q_i^2\right]\,. \tag{3.14}$$

For $x \to 0$ $t(\theta,\nu,x)$ behaves

$$t(\theta,\nu,x) \approx \frac{\sin(\pi\theta)}{\pi}\left(-\ln(x)+\ln(\delta)\right) \tag{3.15}$$

with $\ln(\delta)$ given by

$$\ln(\delta) = 2\psi(1) - \frac{1}{2}\left(\psi(\theta)+\psi(1-\theta)+\psi(\nu)+\psi(1-\nu)\right)\,. \tag{3.16}$$

Thus the dominant pole in the s-channel is

$$\mathscr{A} = i g_s \mathscr{C}\, \mathrm{Tr}\left(\Lambda_{ba}\Lambda_{ab}\Lambda_{bc}\Lambda_{cb}\right)(2\pi)^4 \delta^{(4)}\left(\sum_i^4 k_i\right)\overline{\psi}\cdot\overline{\chi}\,\psi\cdot\chi$$
$$\times \frac{\alpha'^{\frac{3}{2}}}{L_{b^1}L_{b^2}L_{b^3}}\int_0^{0+\varepsilon} dx\, x^{-1+s}\prod_{i=1}^{3}\sum_{p_i,q_i}\left(\frac{x}{\delta}\right)^{\frac{\alpha'}{L^2_{b_i}}p_i^2+\frac{R^2_{x_i}R^2_{y_i}}{\alpha' L^2_{b_i}}q_i^2} . \tag{3.17}$$

For $p_i = q_i = 0$ the amplitude factorizes on the exchange of gauge bosons

$$A_4(k_1,k_2,k_3,k_4) = \mathrm{i}\int \frac{\mathrm{d}^4k\,\mathrm{d}^4k'}{(2\pi)^4}\,\frac{\sum_g A^g_\mu(k_1,k_2,k)A^{g,\mu}(k_3,k_4,k')\delta^{(4)}(k-k')}{k^2-\mathrm{i}\varepsilon} . \tag{3.18}$$

Knowing the form of the three point amplitude allows us to normalize the amplitude. In eq (3.18) we sum over all polarizations (vector index μ) and all colors (adjoint index g) that can be exchanged. The three-point amplitude describing the coupling of two fermions to a gauge boson is given by [48]

$$A^g_\mu(k_1,k_2,k_3) = \mathrm{i}\, g_{D6_b}\,(2\pi)^4\delta^{(4)}\left(\sum_{i=1}^3 k_i\right)\overline{\psi}\sigma^\mu\psi\; Tr(\Lambda_{ba}\Lambda_{ab}\Lambda_{bb}) . \tag{3.19}$$

Here Λ_{bb} denotes the Chan-Paton matrix of the exchanged gauge boson and the gauge coupling reads [40] $g^2_{D6_b} = (2\pi)^4\alpha'^{3/2}g_s/\prod_{i=1}^3 2\pi L_{b_i}$. Performing the integral (3.18) and comparing with (3.17) gives for the normalization $\mathscr{C} = 2\pi$, where we used the usual normalization $Tr(\lambda_a\lambda_b) = \frac{1}{2}\delta_{ab}$.

Non-vanishing p_i and q_i in (3.17) indicate exchanges of KK and winding states, respectively. The exchanges of these states probe the geometry of the D-brane configuration and thus are very model-dependent. On the other hand there are higher order poles not originating from the world-sheet instanton contributions that are related to stringy excitations. Including sub-dominant terms of the hypergeometric functions in the limit $x \to 0$ gives

$$\mathscr{A} = 2i\pi g_s\, \mathrm{Tr}\left(\Lambda_{ba}\Lambda_{ab}\Lambda_{bc}\Lambda_{cb}\right)(2\pi)^4\delta^{(4)}\left(\sum_i^4 k_i\right)\overline{\psi}\cdot\overline{\chi}\,\psi\cdot\chi \tag{3.20}$$
$$\times\frac{\alpha'^{\frac{3}{2}}}{L_{b^1}L_{b^2}L_{b^3}}\int_0^{0+\varepsilon} dx\, x^{-1+s}(1+c_1x+c_2x^2+\ldots)\prod_{i=1}^{3}\sum_{p_i,q_i}\left(\frac{x}{\delta}\right)^{\frac{\alpha'}{L^2_{b_i}}p_i^2+\frac{R^2_{x_i}R^2_{y_i}}{\alpha' L^2_{b_i}}q_i^2} .$$

where c_i are angle dependent coefficients. Note that the sub-dominant poles are integer modded indicating that the mass of the exchanged particles is of order M_s, and can be potentially observed at LHC if the string scale is in the TeV range [6, 8]. Theses signals are very similar to the ones observed in the scattering of multiple gauge bosons onto at most two fermions which have been investigated in [23, 25, 26, 49].

t-channel – exchange of light stringy states

In this channel we expect the exchange of a massless scalar in case of preserved supersymmetry as well as additional massive states whose mass is basically given by the product of the intersection angle and the string scale M_s^2. If the intersection angle is small these will be long-lived resonances which in case of a low string scale could be observed at LHC. In addition to these light-stringy excitations one can also observe exchanges of massive stringy states that even in the limit of a vanishing intersection angle remain massive.

In order to perform this analysis we have to determine the behaviour of $I(\theta,\nu,x)$ and $t(\theta,\nu,x)$ in the limit $x\rightarrow 1$. Using the properties of the hypergeometric functions one obtains for $I(\theta,\nu,x)$

$$\lim_{x\rightarrow 1}\frac{1}{2\pi}I(\theta,\nu,x)\sim\Gamma_{1-\theta,\nu,1+\theta-\nu}\,(1-x)^{\theta-\nu}+\Gamma_{\theta,1-\nu,1-\theta+\nu}\,(1-x)^{\nu-\theta}$$

where we define $\Gamma_{\alpha,\beta,\gamma}=\frac{\Gamma(\alpha)\Gamma(\beta)\Gamma(\gamma)}{\Gamma(1-\alpha)\Gamma(1-\beta)\Gamma(1-\gamma)}$. For $t(\theta,\nu,x)$ we distinguish among two different scenarios, depending on which angle is larger

$$\lim_{x\rightarrow 1}t(\theta,\nu,x)=\frac{\sin(\pi(\theta-\nu))}{2\sin(\pi\nu)}\qquad\text{for}\qquad\theta>\nu \tag{3.21}$$

$$\lim_{x\rightarrow 1}t(\theta,\nu,x)=\frac{\sin(\pi(\nu-\theta))}{2\sin(\pi\nu)}\qquad\text{for}\qquad\theta<\nu\,. \tag{3.22}$$

As a result the amplitude behaves according to

$$\begin{aligned}\mathscr{A}=&2i\pi g_s\,\mathrm{Tr}\left(\Lambda_{ba}\Lambda_{ab}\Lambda_{bc}\Lambda_{cb}\right)\overline{\psi}\cdot\overline{\chi}\,\psi\cdot\chi(2\pi)^4\delta^{(4)}\left(\sum_i^4 k_i\right)\int_{1-\varepsilon}^1 dx\,(1-x)^{-\frac{3}{2}+k_2\cdot k_3}\\&\times\left[\left(\Gamma_{1-\theta^1_{ab},1-\theta^1_{bc},\theta^1_{ab}+\theta^1_{bc}}(1-x)^{\theta^1_{ab}+\theta^1_{bc}-1}+\Gamma_{\theta^1_{ab},\theta^1_{bc},2-\theta^1_{ab}-\theta^1_{bc}}(1-x)^{1-\theta^1_{ab}-\theta^1_{bc}}\right)\right]^{-\frac{1}{2}}\\&\times\left[\left(\Gamma_{1-\theta^2_{ab},1-\theta^2_{bc},\theta^2_{ab}+\theta^2_{bc}}(1-x)^{\theta^2_{ab}+\theta^2_{bc}-1}+\Gamma_{\theta^2_{ab},\theta^2_{bc},2-\theta^2_{ab}-\theta^2_{bc}}(1-x)^{1-\theta^2_{ab}-\theta^2_{bc}}\right)\right]^{-\frac{1}{2}}\\&\times\left[\left(\Gamma_{-\theta^3_{ab},-\theta^3_{bc},2+\theta^3_{ab}+\theta^3_{bc}}(1-x)^{1+\theta^3_{ab}+\theta^3_{bc}}+\Gamma_{1+\theta^3_{ab},1+\theta^3_{bc},-\theta^3_{ab}-\theta^3_{bc}}(1-x)^{-\theta^3_{ab}-\theta^3_{bc}-1}\right)\right]^{-\frac{1}{2}}\\&\times\prod_{p_i,q_i}\prod_{i=1}^2 e^{-S^3_{cl}(\theta^i_{ab},1-\theta^i_{bc},p_i)}e^{-S^3_{cl}(\theta^i_{ab},1-\theta^i_{bc},q_i)}e^{-S^3_{cl}(1+\theta^3_{ab},-\theta^3_{bc},p_3)}e^{-S^3_{cl}(1+\theta^3_{ab},-\theta^3_{bc},q_3)}\,,\end{aligned}$$

where $e^{-S^3_{cl}(\theta,\nu,p)}$ takes the form [30, 43–47]

$$e^{-S^3_{cl}(\theta,\nu,p_i)}=\exp\left[-\frac{\pi}{4}\frac{\sin(\pi\theta)\sin(\pi\nu)}{|\sin(\pi(\theta-\nu))|}\frac{L_{b_i}}{\alpha'}p_i^2\right]\,. \tag{3.23}$$

To simplify the analysis further let us assume that we are in the large volume limit. Thus all world-sheet instanton contributions from $p_i,q_i\neq 0$ are negligible. Additionally for the sake of concreteness the intersection angles satisfy

$$\theta^1_{ab}+\theta^1_{bc}<1\qquad\theta^2_{ab}+\theta^2_{bc}<1\qquad|\theta^3_{ab}+\theta^3_{bc}|>1. \tag{3.24}$$

With these assumptions we can pull out the dominant pole and get for the amplitude

$$\begin{aligned}\mathscr{A}=&2i\pi g_s\,\mathrm{Tr}\left(\Lambda_{ba}\Lambda_{ab}\Lambda_{bc}\Lambda_{cb}\right)\overline{\psi}\cdot\overline{\chi}\,\psi\cdot\chi(2\pi)^4\delta^{(4)}\left(\sum_i^4 k_i\right)\\&\times\int_{1-\varepsilon}^1 dx\frac{(1-x)^{-1-\frac{1}{2}\sum_I(\theta^I_{ab}+\theta^I_{bc})+k_2\cdot k_3}}{\Gamma^{\frac{1}{2}}_{1-\theta^1_{ab},1-\theta^1_{bc},\theta^1_{ab}+\theta^1_{bc}}\,\Gamma^{\frac{1}{2}}_{1-\theta^2_{ab},1-\theta^2_{bc},\theta^2_{ab}+\theta^2_{bc}}\,\Gamma^{\frac{1}{2}}_{-\theta^3_{ab},-\theta^3_{bc},2+\theta^3_{ab}+\theta^3_{bc}}}\\&\times\left[\left(1+c_1(1-x)^{2(1-\theta^1_{ab}-\theta^1_{bc})}\right)\left(1+c_2(1-x)^{2(1-\theta^2_{ab}-\theta^2_{bc})}\right)\left(1+c_3(1-x)^{2(-\theta^3_{ab}-\theta^3_{bc}-1)}\right)\right]^{-\frac{1}{2}}\,.\end{aligned} \tag{3.25}$$

Here the c_i's are given by

$$c_1 = \frac{\Gamma_{1-\theta^1_{ab},1-\theta^1_{bc},\theta^1_{ab}+\theta^1_{bc}}}{\Gamma_{\theta^1_{ab},\theta^1_{bc},2-\theta^1_{ab}-\theta^1_{bc}}} \qquad c_2 = \frac{\Gamma_{1-\theta^2_{ab},1-\theta^2_{bc},\theta^2_{ab}+\theta^2_{bc}}}{\Gamma_{\theta^2_{ab},\theta^2_{bc},2-\theta^2_{ab}-\theta^2_{bc}}} \qquad c_3 = \frac{\Gamma_{-\theta^3_{ab},-\theta^3_{bc},2+\theta^3_{ab}+\theta^3_{bc}}}{\Gamma_{1+\theta^3_{ab},1+\theta^3_{bc},-\theta^3_{ab}-\theta^3_{bc}}} \,.$$

In the case of preserved supersymmetry ($\sum_I \theta^I_{ab} = \sum_I \theta^I_{bc} = 0$) one indeed observes the exchange of a massless scalar [3]. This particle is identified with ϕ whose vertex operator is displayed in eq. (3.8).

The sub-dominant poles reveal the exchanges of massive scalar exchanges whose mass scales as $M^2 \sim \theta^I_{ca} M^2_s$. Including the subdominant poles we get for $x \to 1$

$$\begin{aligned}&\left[\left(1+c_1(1-x)^{2(1-\theta^1_{ab}-\theta^1_{bc})}\right)\left(1+c_2(1-x)^{2(1-\theta^2_{ab}-\theta^2_{bc})}\right)\left(1+c_3(1-x)^{2(-\theta^3_{ab}-\theta^3_{bc}-1)}\right)\right]^{-\frac{1}{2}}\\ &\quad\simeq 1+c_1(1-x)^{2(1-\theta^1_{ab}-\theta^1_{bc})}+c_2(1-x)^{2(1-\theta^2_{ab}-\theta^2_{bc})}+c_3(1-x)^{2(-\theta^3_{ab}-\theta^3_{bc}-1)}+\ldots\end{aligned}$$

For concreteness we assume that $1-\theta^1_{ab}-\theta^1_{bc} = -\theta^1_{ca}$ is small and positive. Then the amplitude takes the following form

$$\mathscr{A} = \overline{\psi}\cdot\overline{\chi}\,\psi\cdot\chi \int_{1-\varepsilon}^{1} dx\,(1-x)^{-1+k_2\cdot k_3}\, Y^2_{\psi\chi\phi}\left(1+c_1(1-x)^{2(1-\theta^1_{ab}-\theta^1_{bc})}+\ldots\right). \tag{3.26}$$

The first sub-dominant term suggests that there is a particle with mass $M^2 = -2\theta^1_{ca}M^2_s$ exchanged.

As we have discussed in the beginning of this section, the spectrum in the ca sector indeed reveals a particle with small positive mass $-2\theta^1_{ca}M^2_s$, namely the scalar $\widehat{\phi}$, whose vertex operator is given in eq. (3.10). Let us stress that there is no coupling to the lightest massive field $\tilde{\phi}$, which one would have naively expected. This is due to the fact that the two bosonic twist fields σ do not couple to the excited twist field τ, but they only couple to an even excited twist field [43]. In agreement with the latter an inspection of higher poles reveals that the next lightest state exchanged has a mass $-4\theta^1_{ca}M^2_s = 2M^2$.

A detailed analysis of the next-lighter massive states while straight-forward is beyond the scope of the present investigation. Similarly we do not analyse (higher spin) massive states, whose masses do not vanish for small angles, but we expect similar results as derived in [22,23,25,26,49]. Such an analysis would require a more detailed analysis of the sub-dominant poles of the hypergeometric functions. Note that while signals induced by light stringy states at colliders could be rather difficult to recognize and discriminate from other kinds of Physics Beyond the Standard Model, still these signals are expected to be observed first. Moreover, at higher energy scales one eventually will observe higher spin state signatures, which then hint towards a stringy nature.

4. Conclusion

We carefully study the spectrum of open strings localized at the intersections of D6-branes. We give a prescription how to derive to any excitation its corresponding vertex operator. One has to pay particular attention to the signs of the intersection angles [23,30,45,50] since the relevant twist fields depend crucially on those.

[3] In the non-SUSY case the lightest exchange particle has mass $M^2 = \frac{1}{2}\sum_{I=1}^{3}\left(\theta^I_{ab}+\theta^I_{bc}\right)$.

We argue that the masses of the lightest states scale as $M_\theta^2 \approx \theta M_s^2$ and can thus be parametrically smaller than the string scale if the relevant angle is small. Furthermore we considerer processes that can expose these light stringy states in their intermediate channel. Relying on previous analysis, we have computed 4-point scattering amplitudes of 'twisted' open strings and studied their factorization in the s- and t-channel confirming the presence of the sought for states as subdominant poles in the latter. We have found that only evenly excited 'twisted' open strings are exchanged in the t-channel, quite differently from what happens for the parent closed-string amplitudes.

We have not analysed in any detail the poles corresponding to massive, possibly higher spin, states which remain massive even when some angles are small. Their analysis is tedious and presents significant analogies with the analysis in [23, 25, 26, 49]. Assuming a scenario with large extra dimensions and a low scale string tension proves to be realized in Nature, the spectrum of string excitations may be rather 'irregular' or at least look very different to the regularly spaced Regge recurrences of the good old Veneziano model. Signals at colliders could be rather difficult to recognize and discriminate from other kinds of Physics Beyond the Standard Model. Yet, the possibility that the lightest massive string excitations be just behind the corner makes worth sharpening our predictions and/or generalizing it to phenomenologically more viable models, possibly including the effect of closed string fluxes and non-perturbative effects.

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