PROCEEDINGS OF SCIENCE

# Yukawa's of light stringy states (at D-brane intersections)

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In [1], we use 4-point amplitudes initially to normalise the vertex operators for massless and massive strings living at D-brane intersections and next to evaluate Yukawa couplings between these states. Our results can be used in order to evaluate physical processes which might be visible at LHC.

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## 1. Introduction

D-brane string vacua has been proven to be a very successful framework for semi-realistic model building[1]. Open strings stretched between stacks of some lower-dimensional hyperplanes, called D-branes, describe *gauge fields* (when the two endpoints end on the same stack of branes) and *chiral matter* (when the endpoints end on different stacks). Therefore, all Standard Model fields can be described by open strings while gravity is mediated by closed strings [2].

A key feature of these vacua is that the string scale can be very low, even at a few TeV range [26–28], and stringy effects become possible candidates for physics beyond the Standard Model (see for example anomalous $Z'$ [29–42], Kaluza Klein states [43–49], or purely stringy signatures [50–63][3]).

Another important prediction of all D-brane realizations of the Standard Model is that each matter field is followed by a whole tower of massive copies. That relays on the fact that a string living at the intersection of two different stacks of D-branes, vibrates with frequencies proportional to the angle between the branes $\theta$ generating a whole tower of massive modes with the same quantum numbers (realised by the endpoints). Therefore, the lowest modes on the intersections realise the electrons, the quarks, the Higgs etc and the highest modes as massive copies of them.

If the string scale $M_s$ is at a few TeV range and some angles between the branes are very small, these copies can be very light (*light stringy states*) and they might be visible in future experiments [66–69].

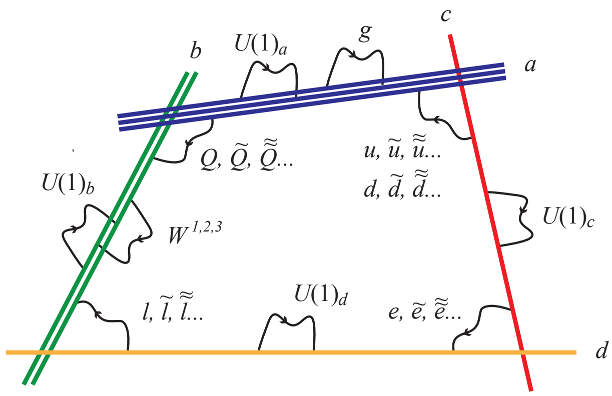


**Figure 1:** Local D-brane realization of the Standard Model. Each field is accompanied by a whole tower of massive excitations.

Even though the tower of massive stringy states looks like KK towers, there are differences. First of all the mass gaps at KK models are typically universal. In intersecting D-brane models each tower of states lives on different intersections with unique angle and therefore unique mass gap. Second the mass gaps in KK modes are typically proportional to the moding ($\sim m/R$) whereas in intersecting D-brane models they are proportional to the square root of the moding ($\sim \sqrt{m}M_s$).

[1] For original work, see [2–14] and for recent reviews on D-brane model building, see [15–19] and references therein.

[2] For original work on local D-brane configurations, see [20, 21]. For a systematic analysis of local D-brane configurations, see [22–25].

[3] For recent reviews, see [64, 65].

Finally, interactions which are prohibited due to conservation of internal momentum (for example the decay of an excited KK mode to two unexcited modes) are allowed in intersecting D-brane models is different from zero [1].

In [1] we compute Yukawa couplings between different stringy states of the towers living at D-brane intersections. We first evaluate 4-point amplitudes and by factorization we first normalise our relevant vertex operators (VO) [70, 71] and next we compute the Yukawa couplings. In this proceeding we sketch the procedure giving some key point of the evaluation. More details can be found in the original work [1].

The plan is as follows. In Section 2, we describe the set-up of intersecting D6-branes on tori and present the expressions for both massless and massive BRST invariant vertex operators. In Section 3 we compute scattering amplitudes on the disk (tree level). Normalization problems are solved by first considering amplitudes that expose vector boson exchange on one or both channels ($s$ and $t$).

## 2. Setup and vertex operators

Consider three stacks of D6-branes wrapping factorizable 3-cycles on a six-torus $T^6 = T^2 \times T^2 \times T^2$, labeled by $a$, $b$ and $c$. In each torus $T_I^2$ ($I = 1,2,3$) two stacks $a$ and $b$ are intersecting at an angle $\theta_{ab}^I = \pi a_{ab}^I$ (next we will use the angles $a_{ab}^I$ subtracting $\pi$ for simplicity). Supersymmetry requires $\pm a_{ab}^1 \pm a_{ab}^2 \pm a_{ab}^3 = 0 \mod 2$, for some choice of signs[4]. For non-vanishing Yukawa couplings, we take

$$a_{ab}^1 + a_{ab}^2 + a_{ab}^3 = 0 \qquad 0 < a_{ab}^1 < 1 \qquad 0 < a_{bc}^1 < 1 \qquad -1 < a_{ca}^1 < 0 \tag{2.1}$$

$$a_{bc}^1 + a_{bc}^2 + a_{bc}^3 = 0 \qquad 0 < a_{ab}^2 < 1 \qquad 0 < a_{bc}^2 < 1 \qquad -1 < a_{ca}^2 < 0 \tag{2.2}$$

$$a_{ca}^1 + a_{ca}^2 + a_{ca}^3 = -2 \quad -1 < a_{ab}^3 < 0 \quad -1 < a_{bc}^3 < 0 \quad -2 < a_{ca}^3 < 0 \tag{2.3}$$

Strings ending on two different stacks of branes give rise to a massless (chiral) spectrum with multiplicity given by the number of intersections and a massive spectrum. The VO's for the lowest string modes (massless and first massive) living at intersections are given by [70]

$$V^{(-1)}_{\phi_0=\phi_0^{ab}} = C_{\phi_0} e^{-\phi_{10}} \phi_0 e^{-\varphi} \sigma_{a_{a,b}^1} \sigma_{a_{a,b}^2} \sigma_{1+a_{a,b}^3} e^{i[a_{a,b}^1\varphi_1 + a_{a,b}^2\varphi_2 + (a_{a,b}^3+1)\varphi_3]} e^{ikX} \tag{2.4}$$

$$V^{(-\frac{1}{2})}_{\psi_0=\chi_0^{bc}} = C_{\psi_0} e^{-\phi_{10}} \psi_0^\alpha S_\alpha e^{-\frac{\varphi}{2}} \sigma_{a_{b,c}^1} \sigma_{a_{b,c}^2} \sigma_{1+a_{b,c}^3} e^{i[(a_{b,c}^1-\frac{1}{2})\varphi_1 + (a_{b,c}^2-\frac{1}{2})\varphi_2 + (a_{b,c}^3+\frac{1}{2})\varphi_3]} e^{ikX} \tag{2.5}$$

$$V^{(-\frac{1}{2})}_{\chi_0=\chi_0^{ca}} = C_{\chi_0} e^{-\phi_{10}} \chi_0^\alpha S_\alpha e^{-\frac{\varphi}{2}} \sigma_{1+a_{c,a}^1} \sigma_{1+a_{c,a}^2} \sigma_{1+a_{c,a}^3} e^{i[(a_{c,a}^1+\frac{1}{2})\varphi_1 + (a_{c,a}^2+\frac{1}{2})\varphi_2 + (a_{c,a}^3+\frac{1}{2})\varphi_3]} e^{ikX} \tag{2.6}$$

$$\begin{aligned} V^{(-\frac{1}{2})}_{\psi_1=\chi_1^{bc}} &= C_{\psi_1} e^{-\phi_{10}} \psi_1^\alpha S_\alpha e^{-\frac{\varphi}{2}} \tau_{a_{b,c}^1} \sigma_{a_{b,c}^2} \sigma_{1+a_{b,c}^3} e^{i[(a_{b,c}^1-\frac{1}{2})\varphi_1 + (a_{b,c}^2-\frac{1}{2})\varphi_2 + (a_{b,c}^3+\frac{1}{2})\varphi_3]} e^{ikX} \\ &+ C_{\tilde\psi_1} e^{-\phi_{10}} \tilde\psi^\dagger_{1\dot\alpha} C^{\dot\alpha} e^{-\frac{\varphi}{2}} \sigma_{a_{b,c}^1} \sigma_{a_{b,c}^2} \sigma_{1+a_{b,c}^3} e^{i[(a_{b,c}^1+\frac{1}{2})\varphi_1 + (a_{b,c}^2-\frac{1}{2})\varphi_2 + (a_{b,c}^3+\frac{1}{2})\varphi_3]} e^{ikX} \end{aligned} \tag{2.7}$$

$$\begin{aligned} V^{(-\frac{1}{2})}_{\chi_1=\chi_1^{ca}} &= C_{\chi_1} e^{-\phi_{10}} \chi_1^\alpha S_\alpha e^{-\frac{\varphi}{2}} \sigma_{1+a_{c,a}^1} \sigma_{1+a_{c,a}^2} \tau_{1+a_{c,a}^3} e^{i[(a_{c,a}^1+\frac{1}{2})\varphi_1 + (a_{c,a}^2+\frac{1}{2})\varphi_2 + (a_{c,a}^3+\frac{1}{2})\varphi_3]} e^{ikX} \\ &+ C_{\tilde\chi_1} e^{-\phi_{10}} \tilde\chi^\dagger_{1\dot\alpha} C^{\dot\alpha} e^{-\frac{\varphi}{2}} \sigma_{1+a_{c,a}^1} \sigma_{1+a_{c,a}^2} \sigma_{1+a_{c,a}^3} e^{i[(a_{c,a}^1+\frac{1}{2})\varphi_1 + (a_{c,a}^2+\frac{1}{2})\varphi_2 + (a_{c,a}^3-\frac{1}{2})\varphi_3]} e^{ikX} \end{aligned} \tag{2.8}$$

[4] Semi-realistic MSSM constructions on factorizable orbifolds can be found in [22, 72–94].

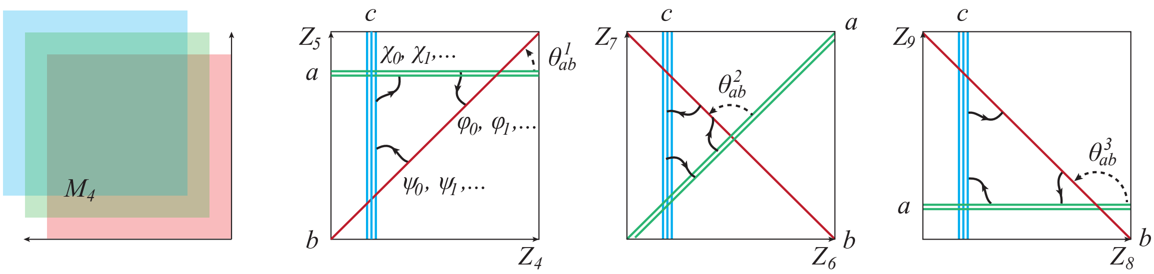


**Figure 2:** A simple configuration of three stacks of D6-branes in a torus $T_I^2$. The angles in the figure are large for illustrative purposes.

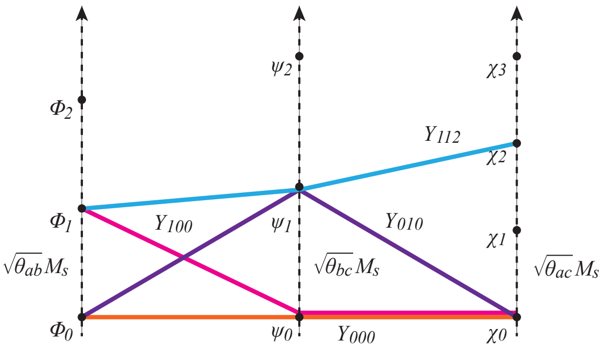

**Figure 3:** We evaluate Yukawas $Y_{ijk}\Phi_i\psi_j\chi_k$ where $i,j,k = 0,1,2...$ the different massive modes, which are denoted by the lines connecting the different modes.

and the masses of the corresponding fields are

$$\textit{massless}\ : \quad m^2_{\phi_0} = 0 \quad , \quad m^2_{\psi_0} = 0 \quad , \quad m^2_{\chi_0} = 0 \tag{2.9}$$

$$\textit{massive}\ : \quad\quad , \quad m^2_{\psi_1} = a^1_{bc}/\alpha' \quad , \quad m^2_{\chi_1} = (1-|a^3_{ca}|)/\alpha' \ . \tag{2.10}$$

Having the expression for the VO's of the massless and massive states living at intersections, we evaluate string amplitudes and consequently the Yukawa couplings.

## 3. Strategy

Three-point amplitudes (Yukawas also) are ambiguous due to the unnormalised VO's of the incoming fields. Therefore, we have to start by four-point amplitudes and by factorisations normalise the corresponding VO's and then evaluate physical amplitudes.

In order to normalise our fields we evaluate their couplings to gauge fields and we equate them with the known values. With normalised VO's at hand we proceed to Yukawas. Schematically the strategy is given in the diagram 4.

### 3.1 The amplitude $\mathscr{A}(\bar\psi_0, \psi_0, \chi_0, \bar\chi_0)$

Following the strategy above we compute specific 4-point amplitudes. Here, we present few examples, and more details can be found at our original work [1].

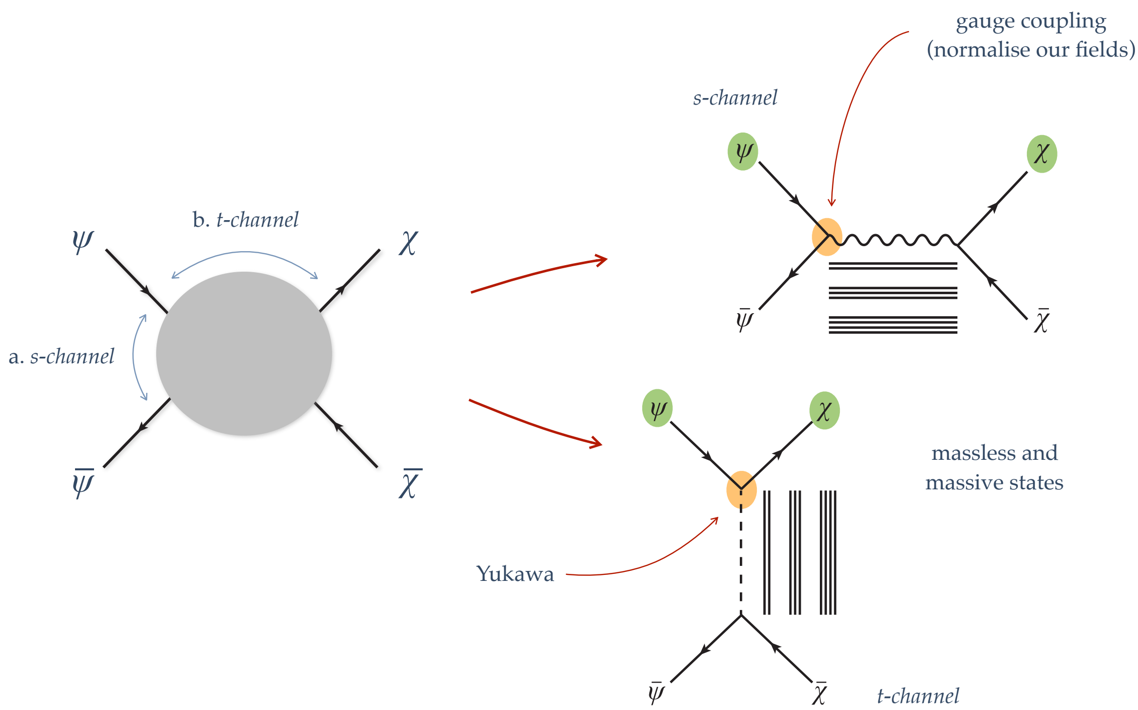


**Figure 4:** Factorization of the 4-point amplitude: a. At the lowest level of the s-channel we have the exchange of a gauge field which helps to normalise our fields. b. At t-channel we extract the various Yukawa couplings.

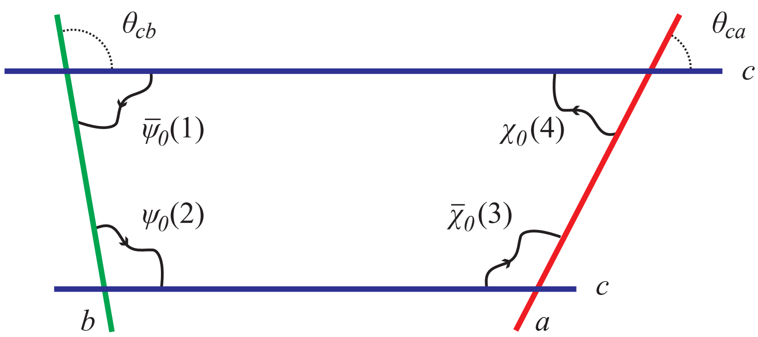


**Figure 5:** A configuration of branes describing the scattering between the spinors $\psi_0$ and $\chi_0$.

The amplitude $\mathscr{A}(\bar\psi_0, \psi_0, \chi_0, \bar\chi_0)$ with two-independent angles yields

$$\begin{aligned}\mathscr{A}(\bar\psi_0^{cb}, \psi_0^{bc}, \chi_0^{ca}, \bar\chi_0^{ac}) &= \langle V_{\bar\psi_0}^{(-\frac{1}{2})} V_{\psi_0}^{(-\frac{1}{2})} V_{\chi_0}^{(-\frac{1}{2})} V_{\bar\chi_0}^{(-\frac{1}{2})}\rangle \\ &= g_{\rm op}^2 \alpha' \psi_0(2)\cdot\chi_0(3)\,\bar\psi_0(1)\cdot\bar\chi_0(4) \int_0^1 dx\, x^{\alpha' s-1}(1-x)^{\alpha' t-1} \times \\ &\times \prod_{I=1}^{3} \frac{4\pi^2 K_I^{c,ab}\alpha'}{L_{a,I}^{1/4} L_{b,I}^{5/4} L_{c,I}^{1/2}} \frac{\sqrt{\alpha'}}{L_{c,I} G_1^{(I)}(x)} \sum_{n_I, m_I} e^{-S_{\rm Ham}^{(I)}(m_I, n_I)}\end{aligned} \tag{3.1}$$

where $L_a, I$ is the brane $a$'s length in torus $I$. $I_{ab}$ and similar are the number or intersections between the branes on the torus. $G_1^{(I)}(x)$ are hypergeometric functions $G_1 = {}_2F_1\,(\alpha, 1-\beta; 1; z)$, $G_2 = {}_2F_1\,(1-\alpha, \beta; 1; z)$, $K_I^{c,ab}$ is an overall normalization constant and $S^{(I)}_{\rm Ham}(m_I, n_I)$ is the classical action in "hamiltonian" form:

$$S^{(I)}_{\rm Ham}(m_I,n_I) = \frac{\pi t_I(x)}{\sin\pi|a^I_{bc}|}\left(\frac{4\pi^2\alpha' m_I^2}{L^2_{c,I}} + \frac{\sin^2\pi|a^I_{bc}|L^2_{b,I}}{4\pi^2\alpha'}\frac{I^2_{ca,I}n_I^2}{\gcd^2(|I_{bc,I}|,|I_{ca,I}|)}\right) - 2\pi i\frac{m_I}{L_{c,I}}f_{\chi\psi,I} \quad (3.2)$$

In the $s$ channel the amplitude shows a gauge boson propagating between parallel branes of type $c$, we can use the factorization in this channel to fix $K_I^{c,ab} = \frac{L_a^{1/4}L_b^{5/4}L_{c,I}^{1/2}}{(2\pi)^2\alpha'}$. The vertex operators' normalizations obtained from the amplitudes are

$$C_{A_i} = \sqrt{2\alpha'}\prod_I\left[\frac{\alpha'}{L^2_{i,I}}\right]^{1/4} \quad (3.3)$$

$$C_{\chi_0^{ij}} = (\alpha')^{1/4}\sqrt{2\alpha'}\prod_I\left[\frac{\alpha'}{L_{i,I}L_{j,I}}\right]^{1/4} \qquad C_{\phi_0^{ij}} = \sqrt{2\alpha'}\prod_I\left[\frac{\alpha'}{L_{i,I}L_{j,I}}\right]^{1/4} \quad (3.4)$$

$$C_{\chi_1^{ij}} = (\alpha')^{1/4}\sqrt{2\alpha'}\prod_I\left[\frac{\alpha'}{L_{i,I}L_{j,I}}\right]^{1/4} \qquad C_{\tilde\chi_1^{ij}} = (\alpha')^{1/4}\sqrt{2\alpha'}\prod_I\left[\frac{\alpha'}{L_{i,I}L_{j,I}}\right]^{1/4} \quad (3.5)$$

The factorization in the $t$ channel exhibits Yukawa couplings. After Poisson resummation we get

$$\mathscr{A}(\bar\psi_0,\psi_0,\chi_0,\bar\chi_0) = g^2_{\rm op}\alpha'\psi_0(2)\cdot\chi_0(3)\bar\psi_0(1)\cdot\bar\chi_0(4)\int_0^1 dx\,x^{\alpha's-1}(1-x)^{\alpha't-3/2}\prod_{I=1}^3\sum_{\tilde m_I,n_I}\frac{e^{-S^{(I)}_{\rm Lagr}(\tilde m_I,n_I)}}{2\pi\sqrt{I_I(x)}} \quad (3.6)$$

where $I_I(x)$ are combinations of hypergeometric functions and $S^{(I)}_{\rm Lagr}(\tilde m_I, n_I)$ is the classical action in "lagrangian" form. The limit $x\to 1$ has two kinds of contributions: one purely quantum, due to the expansion of $I_I(x)$, and one classical, due to the expansion of $t(x)$ in the action. The first three orders correspond to factorizations on the poles for the states $\phi_0$, $\phi_1$ and $\phi_2$. Imposing

$$\mathscr{A}(\bar\psi_0,\psi_0,\chi_0,\bar\chi_0) \xrightarrow{t\to n a^1_{ab}/\alpha'} |Y_{n00}|^2\psi_0(2)\cdot\chi_0(3)\frac{1}{t-na^1_{ab}/\alpha'}\bar\chi_0(4)\cdot\bar\psi_0(1) \quad (3.7)$$

The Yukawa's extracted by factorization are given by

$$|Y_{000}| = g_{\rm op}(2\pi)^{-3/4}[\Gamma_{1-a^1_{ab},1-a^1_{bc},-a^1_{ca}}\Gamma_{1-a^2_{ab},1-a^2_{bc},-a^2_{ca}}\Gamma_{-a^3_{ab},-a^3_{bc},-a^3_{ca}}]^{1/4}\prod_{I=1}^3\exp\left[-\frac{A^{(I)}_{\phi\psi\chi}}{2\pi\alpha'}\right] \quad (3.8)$$

$$|Y_{100}| = \frac{|Y_{000}|}{\sqrt{a^1_{ab}}}[\Gamma_{1-a^1_{ab},1-a^1_{bc},-a^1_{ca}}]^{1/2}\sqrt{\frac{2A^{(1)}_{\phi\psi\chi}}{\pi\alpha'}} \quad (3.9)$$

$$|Y_{200}| = \frac{|Y_{000}|}{\sqrt{2}a^1_{ab}}\Gamma_{1-a^1_{ab},1-a^1_{bc},-a^1_{ca}}\left|\frac{2A^{(1)}_{\phi\psi\chi}}{\pi\alpha'}-1\right| \quad (3.10)$$

where $\Gamma_{a,b,c} = \dfrac{\Gamma(a)\Gamma(b)\Gamma(c)}{\Gamma(1-a)\Gamma(1-b)\Gamma(1-c)}$, and $A^{(I)}_{\phi\psi\chi}$ is the area of the triangle defined by the three points $f_{\psi,I}$, $f_{\chi,I}$ and $f_{\phi,I}$ in the torus $I$ given by $A^{(I)}_{\phi\psi\chi} = \frac{\sin\pi|a^I_{bc}|\sin\pi|a^I_{ca}|}{2\sin\pi|a^I_{ab}|}f^2_{\chi\psi,I}$.

For the rest of the Yukawas $Y_{010}, Y_{001}$ etc we use supersymmetric Ward identities [1] and we get

$$|Y_{010}| = \frac{|Y_{000}|}{\sqrt{a^1_{bc}}}[\Gamma_{1-a^1_{ab},1-a^1_{bc},-a^1_{ca}}]^{1/2}\sqrt{\frac{2A^{(1)}_{\phi\psi\chi}}{\pi\alpha'}}\ , \quad |Y_{001}| = \frac{|Y_{000}|}{\sqrt{1+a^3_{ca}}}[\Gamma_{-a^3_{ab},-a^3_{bc},-a^3_{ca}}]^{1/2}\sqrt{\frac{2A^{(3)}_{\phi\psi\chi}}{\pi\alpha'}} \tag{3.11}$$

$$|Y_{020}| = \frac{|Y_{000}|}{\sqrt{2}a^1_{bc}}\Gamma_{1-a^1_{ab},1-a^1_{bc},-a^1_{ca}}\left|\frac{2A^{(1)}_{\phi\psi\chi}}{\pi\alpha'}-1\right|\ , \quad |Y_{002}| = \frac{|Y_{000}|}{\sqrt{2}(1+a^3_{ca})}\Gamma_{-a^3_{ab},-a^3_{bc},-a^3_{ca}}\left|\frac{2A^{(3)}_{\phi\psi\chi}}{\pi\alpha'}-1\right| \tag{3.12}$$

which completes the list of Yukawas involving one massive (1st and 2nd excitations) and two massless states.

### 3.2 Yukawa couplings from amplitudes with massive external legs

For Yukawas of more than one massive particle involved we have to start by 4-pt amplitudes with massive external legs. Following our strategy again, we fix the normalizations of the vertex operators from amplitudes that include gauge bosons, and we evaluate Yukawa couplings with more than one massive field.

Without loss of generality, we take the amplitude $\mathscr{A}(\bar\psi_1, \psi_0, \chi_0, \bar\chi_0)$ with one massive state $\bar\psi_1$ and three massless $\psi_0, \chi_0, \bar\chi_0$:

$$\begin{aligned}\mathscr{A}(\bar\psi_1, \psi_0, \chi_0, \bar\chi_0) = {} & \frac{g^2_{\rm op}}{\sqrt{a^1_{bc}}}\alpha'\psi_0(2)\cdot\chi_0(3)\bar\psi_1(1)\cdot\bar\chi_0(4) \\ & \times\int_0^1 dx\, x^{\alpha's-1}(1-x)^{\alpha't-1}\prod_{I=1}^{3}\frac{2\pi\sqrt{\alpha'}}{L_{c,I}G^{(I)}_1(x)}\sum_{m_I}\frac{2\pi\sqrt{\alpha'}m_1}{\sqrt{I_1(x)}L_{c,I}}e^{-S^{(I)}_{\rm Ham}(m_I,n_I)}\end{aligned} \tag{3.13}$$

Similar to the amplitudes with massless external states, we perform a Poisson resummation over the indices $m_I$ and obtain

$$\begin{aligned}\mathscr{A}(\bar\psi_1, \psi_0, \chi_0, \bar\chi_0) = {} & \frac{g^2_{\rm op}}{\sqrt{a^1_{bc}}}\alpha'\psi_0(2)\cdot\chi_0(3)\bar\psi_1(1)\cdot\bar\chi_0(4)\int_0^1 dx\, x^{\alpha's-1}(1-x)^{\alpha't-3/2-a^1_{ab}} \\ & \times\frac{G^{(1)}_1(x)}{I_1(x)}\prod_{I=1}^{3}\frac{1}{\sqrt{2\pi I_I(x)}}\sum_{\tilde m_I,n_I}\frac{\tilde m_1 L_{c,1}+f_{\chi\psi,1}}{2\pi\sqrt{\alpha'}}e^{-S^{(I)}_{\rm Lagr}(\tilde m_I,n_I)}\end{aligned} \tag{3.14}$$

In the limit $x \to 1$ the leading term is the massless pole due to the chiral exchange in the $(a,b)$ sector. We have fixed all the normalizations that appear in the amplitudes and we already know the Yukawa's $Y^*_{010}Y_{000}$, thus factorization in this channel can be used to check that normalizations are consistent. The subleading terms determine the Yukawa $Y_{110}$

$$|Y_{110}| = |Y_{000}|\left|\frac{2A^{(1)}_{\phi\psi\chi}}{\pi\alpha'}-1\right|\frac{1}{\sqrt{a^1_{ab}a^1_{bc}}}\Gamma_{1-a^1_{ab},1-a^1_{bc},-a^1_{ca}} \tag{3.15}$$

We can obtain the Yukawas $Y_{101}$ and $Y_{011}$ from two amplitudes with massive bosons in the external states. As for the massless case, we can use Ward identities to relate these to amplitudes with only fermions. The explicit expressions of the amplitudes read

$$\mathscr{A}(\bar{\chi}_1^{ac},\chi_0^{ca},\phi_0^{ab},\bar{\phi}_0^{ba}) = \frac{g_{\rm op}^2}{\sqrt{1+a_{ca}^3}}\alpha'\bar{\chi}_1^{ac}(1)\not{k}_4\chi_0^{ca}(2)\int_0^1 dx\,x^{\alpha's-1}(1-x)^{\alpha't-5/2-a_{bc}^3}\times \\ \times\frac{G_1^{(3)}(x)}{I_3(x)}\prod_{I=1}^3\frac{1}{\sqrt{2\pi I_I(x)}}\sum_{\tilde{m}_I,n_I}\frac{\tilde{m}_3L_{a,3}+f_{\psi\phi,3}}{2\pi\sqrt{\alpha'}}e^{-S_{\rm Lagr}^{(I)}(\tilde{m}_I,n_I)} \tag{3.16}$$

$$\mathscr{A}(\bar{\phi}_1,\phi_0,\chi_0,\bar{\chi}_0) = \frac{g_{\rm op}^2}{\sqrt{a_{ab}^1}}\alpha'\bar{\chi}_0^{cb}(4)\not{k}_2\bar{\chi}_0^{bc}(3)\int_0^1 dx\,x^{\alpha's-1}(1-x)^{\alpha't-5/2-a_{ca}^1}\times \\ \times\frac{G_1^{(1)}(x)}{I_1(x)}\prod_{I=1}^3\frac{1}{\sqrt{2\pi I_I(x)}}\sum_{\tilde{m}_I,n_I}\frac{\tilde{m}_1L_{b,1}+f_{\chi\phi,1}}{2\pi\sqrt{\alpha'}}e^{-S_{\rm Lagr}^{(I)}(\tilde{m}_I,n_I)} \tag{3.17}$$

The limit $x\to 1$ yields two new Yukawa's

$$|Y_{011}| = \frac{|Y_{000}|}{\sqrt{a_{bc}^1(1+a_{ca}^3)}}[\Gamma_{1-a_{ab}^1,1-a_{bc}^1,-a_{ca}^1}\Gamma_{-a_{ab}^3,-a_{bc}^3,-a_{ca}^3}]^{1/2}\sqrt{\frac{2A_{\phi\psi\chi}^{(1)}}{\pi\alpha'}\frac{2A_{\phi\psi\chi}^{(3)}}{\pi\alpha'}} \tag{3.18}$$

$$|Y_{101}| = \frac{|Y_{000}|}{\sqrt{a_{ab}^1(1+a_{ca}^3)}}[\Gamma_{1-a_{ab}^1,1-a_{bc}^1,-a_{ca}^1}\Gamma_{-a_{ab}^3,-a_{bc}^3,-a_{ca}^3}]^{1/2}\sqrt{\frac{2A_{\phi\psi\chi}^{(1)}}{\pi\alpha'}\frac{2A_{\phi\psi\chi}^{(3)}}{\pi\alpha'}} \tag{3.19}$$

The amplitude $\mathscr{A}(\bar{\psi}_1,\psi_0,\chi_0,\bar{\chi}_1)$ allows us to determine the Yukawa's $Y_{111}$ and $Y_{211}$.

$$\mathscr{A}(\bar{\psi}_1,\psi_0,\chi_0,\bar{\chi}_1) = \frac{g_{\rm op}^2}{\sqrt{a_{bc}^1(1+a_{ca}^3)}}\alpha'\psi_0(2)\cdot\chi_0(3)\bar{\psi}_1(1)\cdot\bar{\chi}_1(4)\int_0^1 dx\,x^{\alpha's-1}(1-x)^{\alpha't-1} \\ \times\prod_{I=1}^3\frac{4\pi^2\sqrt{\alpha'}}{L_{c,I}G_1^{(I)}(x)}\sum_{m_I,n_I}\frac{\alpha'm_1m_3}{\sqrt{I_1(x)I_3(x)}L_{c,1}L_{c,3}}e^{-S_{\rm Ham}^{(I)}(m_I,n_I)} \tag{3.20}$$

Once again the amplitude does not expose gauge boson exchange, since the sum over the lattice forbids it. To study the $t$ channel we perform the usual Poisson resummation

$$\mathscr{A}(\bar{\psi}_1,\psi_0,\chi_0,\bar{\chi}_1) = \frac{\alpha'g_{\rm op}^2}{\sqrt{a_{bc}^1(1+a_{ca}^3)}}\psi_0(2)\cdot\chi_0(3)\bar{\psi}_1(1)\cdot\bar{\chi}_0(4)\int_0^1 dx\,x^{\alpha's-1}(1-x)^{\alpha't-5/2+a_{ab}^2} \\ \times\frac{G_1^{(1)}(x)G_1^{(3)}(x)}{I_1(x)I_3(x)}\prod_{I=1}^3\frac{1}{\sqrt{2\pi I_I(x)}}\sum_{\tilde{m}_I,n_I}\frac{(\tilde{m}_1L_{c,1}+f_{\psi\chi,1})(\tilde{m}_3L_{c,3}+f_{\psi\chi,3})}{4\pi^2\alpha'}e^{-S_{\rm Lagr}^{(I)}(\tilde{m}_I,n_I)} \tag{3.21}$$

Factorizing this amplitude we obtain the Yukawa's

$$|Y_{111}| = \frac{|Y_{000}|}{\sqrt{a^1_{ab}a^1_{bc}(1+a^3_{ca})}}\Gamma_{1-a^1_{ab},1-a^1_{bc},-a^1_{ca}}\Gamma^{1/2}_{-a^3_{ab},-a^3_{bc},-a^3_{ca}}\left|\frac{2A^{(1)}_{\phi\psi\chi}}{\pi\alpha'}-1\right|\sqrt{\frac{2A^{(3)}_{\phi\psi\chi}}{\pi\alpha'}} \quad (3.22)$$

$$|Y_{211}| = \frac{|Y_{000}|}{\sqrt{2}a^1_{ab}\sqrt{a^1_{bc}(1+a^3_{ca})}}\Gamma^{3/2}_{1-a^1_{ab},1-a^1_{bc},-a^1_{ca}}\Gamma^{1/2}_{-a^3_{ab},-a^3_{bc},-a^3_{ca}}\left|\frac{2A^{(1)}_{\phi\psi\chi}}{\pi\alpha'}-3\right|\sqrt{\frac{2A^{(1)}_{\phi\psi\chi}}{\pi\alpha'}\frac{2A^{(3)}_{\phi\psi\chi}}{\pi\alpha'}} \quad (3.23)$$

## 4. Conclusions

In [1] we have evaluated the Yukawa couplings between light stringy states and SM fields. Our results can be used in order to built an effective field theory including SM fields and light stringy states and study specific decays. If the string scale is at the few TeV range, decays of such particles might be the first stringy effect which might be visible at LHC or future experiments.

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