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Light stringy state production

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This proceeding is based on [1], where we study amplitudes involving two gauge bosons and two light stringy states. The latter states correspond to the lightest stringy excitations localized at the intersection of two D-brane stacks whose masses can be parametrically smaller than the string scale. Given a low string scale scenario with a string scale of a few TeV such states might be observable at the LHC.

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1. Introduction

D-brane model building has been very successful in giving semi-realistic descriptions of the Standard Model (SM) and providing a rich framework for physics beyond the SM[1]. In contrast to heterotic model building D-brane compactifications allow for a low string scale M_s, even of just a few TeV [6–8] [2] providing a solution to the hierarchy and cosmological constant problems but also leading to interesting signatures which could be observed at the LHC such as signs of anomalous Z' physics, Kaluza Klein states and Regge stringy states (see, e.g. [12–49]).

In these D-brane compactifications the gauge bosons are strings attached to a stack of lower-dimensional hyperplanes, so called D-branes, while the chiral matter, such as the SM fermions, are strings that appear at the intersection of different D-brane stacks.

The motivation for this work arises from the fact that at the intersection of two D-brane stacks there exists beyond the massless fermion a whole tower of stringy excitations whose mass scales as $M^2 \sim \theta M_s^2$, where θ denotes the intersection angle between these two D-brane stacks [50].

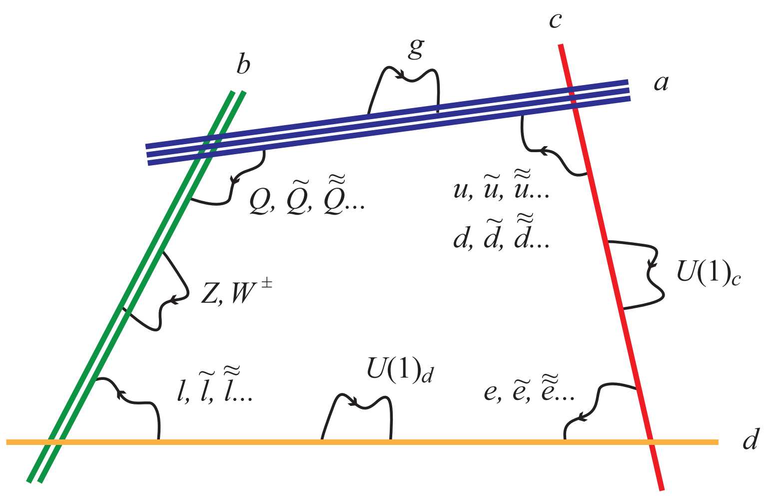

Figure 1: Local D-brane realization of the Standard Model.

In figure 1 we depict a D-brane SM realization. While the SM gauge bosons live on the world volume of the D-brane stacks the SM fermions are localized at the intersections of different D-brane stacks. At each intersection there exist a tower of massive stringy excitations that can be significantly lighter than the string scale for some regions in the parameter space (aka the *light stringy states*). Let us stress that in this setup one has for each massless fermion a separate tower of stringy excitations with a different mass spacing. In that respect it can be easily discriminated from Kaluza-Klein scenarios, in which generically the mass gap is for each particle the same.

In [1] we investigate properties of those light stringy states. More precisely, we compute 4-point tree level string scattering amplitudes containing two gauge bosons and two light stringy states. A crucial ingredient in this calculation is the exact knowledge of the vertex operators of

[1]For recent reviews on D-brane model building as well as specific MSSM D-brane constructions, see [2–5] and references therein.

[2]The lower bound of the string scale is at 4.5 TeV [9–11].

the light stringy states, which we derive by applying the dictionary laid out in [50]. With the corresponding vertex operators at hand we compute the full string scattering amplitude and sum over all polarization and gauge configurations of the squared amplitude. Finally we apply our results to a D-brane realization of the Standard Model, thus bringing it into a form suitable to comparison to observed LHC data.

2. Vertex operators

For concreteness we perform the computation in the type IIA setting, however our results can be easily translated to the T-dual type IIB configuration with D-brane carrying magnetic fluxes. The basic building blocks of type IIA constructions are D6-branes. The SM gauge group is then realized by strings with both ends on the same stack of D6-branes. On the other hand the chiral SM matter fields are localized at the intersections of two such D6-brane stacks.

In order to ensure calculability, we assume that the SM is located at an area of the Calabi-Yau manifold which looks like a factorizable six-torus $T^6 = T^2 \times T^2 \times T^2$ [3] and D6-branes wrap 3 cycles within this local factorizable six-torus (see [30] for more details). More precisely, the D6-branes wrap one-cycles in each of the two tori. Thus the intersection of two D6-brane stacks is decried by three intersection angles, namely by θ^1_{ab}, θ^2_{ab} and θ^3_{ab}. Even though it will not be crucial for our analysis let us display the constraint on the three intersection angles to preserve $\mathcal{N} = 1$ supersymmetry

$$\theta^1_{ab} + \theta^2_{ab} + \theta^3_{ab} = 0\,. \tag{2.1}$$

Here the angles are multiples of π. Without loss of generality we can assume that (2.1) is satisfied with the choice

$$\theta^1_{ab} > 0 \qquad \theta^2_{ab} > 0 \qquad \theta^3_{ab} < 0\,. \tag{2.2}$$

In figure 2 we depict the massless fermionic mode and their complex conjugate counterpart $(\psi, \overline{\psi})$ living at the intersections of two D6-branes. In addition we display the lightest stringy excitation $(\widetilde{\psi}, \widetilde{\overline{\psi}})$ which have the same quantum numbers as $(\psi, \overline{\psi})$ however are massive.

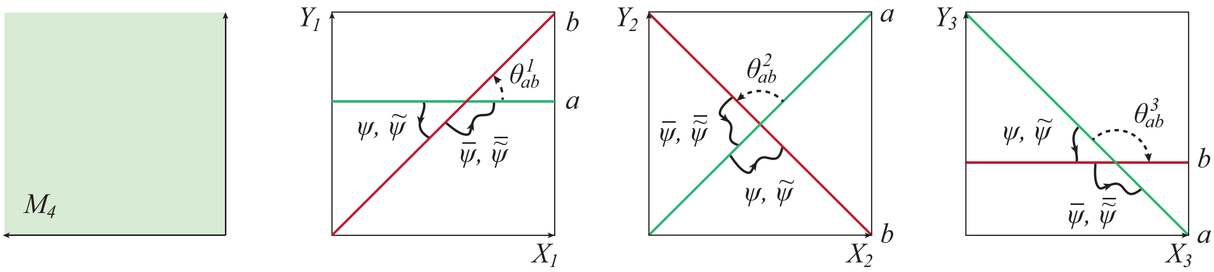

Figure 2: Two intersecting D-brane stacks on a $T^2 \times T^2 \times T^2$.

[3]For semi-realistic MSSM constructions compactified on a factorizable orbifold, see [51–73].

The aim of this work is to compute the string four-point scattering amplitude of two gauge bosons into the first stringy excitation of the massless chiral fermion localized at the intersection of two D6-branes. More precisely, we want to calculate the string amplitude

$$\mathscr{M} = \int \frac{\prod_{i=1}^{4} dz_i}{V_{CKG}} \left\langle V_A^{(0)}[z_1,\varepsilon_1,k_1]\, V_A^{(-1)}[z_2,\varepsilon_2,k_2]\, V_{\widetilde{\psi}}^{(-1/2)}[z_3,v_3,\overline{u}_3,k_3]\, V_{\overline{\widetilde{\psi}}}^{(-1/2)}[z_4,\overline{v}_4,u_4,k_4] \right\rangle , \tag{2.3}$$

where A denotes the gauge field living on the D-brane.

The corresponding vertex operators (VO's) in (2.3) are listed below:

- The VO's of a gauge boson in the (-1) and (0)-ghost picture take the form

$$V_A^{(-1)}[z,\varepsilon,k] = g_A [T^a]^{\alpha_1}_{\alpha_2}\, e^{-\phi}\, \varepsilon^\mu\, \psi_\mu e^{ikX} \tag{2.4}$$

$$V_A^{(0)}[z,\varepsilon,k] = \frac{g_A}{\sqrt{2\alpha'}} [T^a]^{\alpha_1}_{\alpha_2}\, \varepsilon^\mu \left\{ \partial X_\mu - 2i\alpha'(k\cdot\psi)\,\psi_\mu \right\} e^{ikX} \ . \tag{2.5}$$

  Here ε^μ is the polarization vector, while $[T^a]^{\alpha_1}_{\alpha_2}$ denotes the Chan-Paton factor, where a denotes the D-brane stack on which the gauge boson is localized and the indices α_1 and α_2 describe the two string ends. Finally, the string vertex coupling g_A can be determined by computing the three gauge boson disk scattering amplitude [4]

$$g_A = \sqrt{2\alpha'}\, g_{Dp_a} \ . \tag{2.6}$$

- Before we discuss the VO's of the massive fields that appear in (2.3), we present the massless fermionic modes localized at the intersection of two D6 brane stacks a and b:

  – The fermionic massless mode which is localized at such intersection is the R-vacuum $|\,\theta\,\rangle_R$, with vertex operator for the concrete choice of intersection angles (2.2) [5]

$$|\,\theta\,\rangle_R^{ab}: \quad V_\psi^{(-1/2)} = g_\psi\, [T^{ab}]^{\beta_1}_{\alpha_1}\, e^{-\varphi/2}\, v^\alpha\, S_\alpha \prod_{I=1}^{2} \sigma_{\theta^I_{ab}}\, e^{i\left(\theta^I_{ab}-\frac{1}{2}\right)H_I}\, \sigma_{1+\theta^3_{ab}}\, e^{i\left(\theta^3_{ab}+\frac{1}{2}\right)H_3}\, e^{ikX} \ . \tag{2.7}$$

  The Chan-Paton factors indicate that the string is stretching between the two D-brane stacks a and b. The indices α_1 and β_1 run from one to the dimension of the fundamental representation of the gauge group living on the respective D-brane stack a and b. The internal part of the vertex operators contains bosonic and fermionic twist fields σ_{α_I} and $e^{-i\alpha_I H_I}$ introduced to account for the mixed boundary conditions of the open string stretched between the two intersecting D-brane stacks.

  The GSO projection determines the chirality of the four-dimensional polarization spinor and BRST invariance condition (for more details see [1]) gives:

$$\alpha' k^2 = 0 \qquad\qquad v^\alpha \sqrt{\alpha'} k_\mu \sigma^\mu_{\alpha\dot\alpha} = 0 \ , \tag{2.8}$$

[4] Here we use the normalization factor $C_{D2} = \alpha'^{-2}\, g^{-2}_{Dp_a}$ for disk amplitudes containing strings localized solely on a single D-brane stack [30, 75, 76].

[5] For a detailed discussion on vertex operators of massless states for arbitrary intersection angles, see [77, 78], for a generalization to massive states see [50] and for a discussion on instantonic states at the intersection of D-instanton and D-brane at arbitrary angles, see [79].

where the second part of (2.8) is the equation of motion (EOM) for a massless Weyl fermion.

– Similarly to the above, one obtains for the complex-conjugate massless fermionic mode $\overline{\psi}$

$$|\,1-\theta\,\rangle_R^{ba}: \quad V_{\overline{\psi}}^{(-1/2)} = g_{\overline{\psi}}\,[T^{ba}]_{\beta_1}^{\alpha_1}\,e^{-\varphi/2}\,\overline{v}_{\dot{\alpha}}\,\overline{S}^{\dot{\alpha}}\,\prod_{I=1}^{2}\sigma_{1-\theta_{ab}^I}\,e^{-i\left(\theta_{ab}^I-\frac{1}{2}\right)H_I}\,\sigma_{-\theta_{ab}^3}\,e^{-i\left(\theta_{ab}^3+\frac{1}{2}\right)H_3}\,e^{ikX}\,, \tag{2.9}$$

where GSO projection ensures that it is a spinor with opposite chirality to (2.7) and BRST invariance provides a similar massless equation

$$\alpha' k^2 = 0 \qquad\qquad \overline{v}_{\dot{\alpha}}\sqrt{\alpha'}k^{\mu}\overline{\sigma}_{\mu}^{\dot{\alpha}\alpha} = 0\,. \tag{2.10}$$

Finally, the string vertex couplings g_{ψ}, $g_{\overline{\psi}}$ can be derived by calculating the three point amplitude $\langle A\psi\overline{\psi}\rangle$. One obtains

$$g_{\psi} = g_{\overline{\psi}} = \sqrt{2}\alpha'^{\frac{3}{4}}e^{\phi_{10}/2}\,, \tag{2.11}$$

where we used the normalization factor for a disk scattering amplitude containing strings stretched between different D-brane stacks that is given by $\widetilde{C}_{D2} = \alpha'^{-2}e^{-\phi_{10}}$.

- Let us turn to the VO's of the lightest massive stringy excitations

  – The lightest fermionic massive state localized at the intersection the two D-brane stacks a and b is generated by the bosonic and fermionic creation operators, $\alpha_{-\theta_{ab}^1}$ and $\psi_{-\theta_{ab}^1}$, acting on the R-vacuum $|\,\theta\,\rangle_R$. Those two Weyl spinors combine into a massive Dirac spinor $\widetilde{\psi} = (\widetilde{v}_{\alpha}, \overline{\widetilde{u}}_{\dot{\alpha}})$ where $\widetilde{v}$ and $\overline{\widetilde{u}}$ are created by $\alpha_{-\theta_{ab}^1}$ and $\psi_{-\theta_{ab}^1}$, respectively. Without loss of generality let us assume that θ_{ab}^1 is the smallest of all three intersection angles. Then the mass of this Dirac fermion is $\alpha' m^2 = \theta_{ab}^1$. and its corresponding vertex operator is given by [1]

$$V_{\widetilde{\psi}}^{(-1/2)} = g_{\widetilde{\psi}}\,[T^{ab}]_{\alpha_1}^{\beta_1}e^{-\varphi/2}\left(\frac{\widetilde{v}_{\alpha}}{\sqrt{\theta_{ab}^1}}S^{\alpha}\,\tau_{\theta_{ab}^1}\,e^{i(\theta_{ab}^1-\frac{1}{2})H_1} + \overline{\widetilde{u}}_{\dot{\alpha}}\,S^{\dot{\alpha}}\,\sigma_{\theta_{ab}^1}\,e^{i(\theta_{ab}^1+\frac{1}{2})H_1}\right)$$
$$\times\,\sigma_{\theta_{ab}^2}\,e^{i(\theta_{ab}^2-\frac{1}{2})H_2}\,\sigma_{1+\theta_{ab}^3}\,e^{i(\theta_{ab}^3+\frac{1}{2})H_3}\,e^{ikX}\,. \tag{2.12}$$

The BRST condition requires

$$\alpha' k^2 = -\theta_{ab}^1 \tag{2.13}$$

$$\widetilde{v}_{\alpha}\sqrt{\alpha'}\,k^{\mu}\sigma_{\mu}{}^{\alpha\dot{\alpha}} + \sqrt{\theta_{ab}^1}\,\overline{\widetilde{u}}_{\dot{\alpha}} = 0 \qquad\qquad \overline{\widetilde{u}}_{\dot{\alpha}}\sqrt{\alpha'}\,k^{\mu}\overline{\sigma}_{\mu}^{\dot{\alpha}\alpha} + \sqrt{\theta_{ab}^1}\,\widetilde{v}_{\alpha} = 0\,, \tag{2.14}$$

where the second line (2.14) is the Dirac equation written in terms of the left and right-handed Weyl components of the massive Dirac spinor $\widetilde{\psi}$.

– The complex conjugate $\overline{\widetilde{\psi}}$ of the previous massive fermion is coming from the action of $\widetilde{\alpha}_{-\theta^1_{ab}}$ and $\widetilde{\psi}_{-\theta^1_{ab}}$ on the R-vacuum $|\,1-\theta\,\rangle^{ba}_R$ where the "tilde" denotes the creation operators of a system of intersecting D6-branes with opposite intersection angles[6]. The corresponding vertex operator takes the form

$$V^{(-1/2)}_{\overline{\widetilde{\psi}}} = g_{\overline{\widetilde{\psi}}}\, [T^{ba}]^{\alpha_1}_{\beta_1} e^{-\varphi/2} \left(\frac{\overline{\widetilde{v}}_{\dot\alpha}}{\sqrt{\theta^1_{ab}}} S^{\dot\alpha}\, \widetilde{\tau}_{1-\theta^1_{ab}}\, e^{-i(\theta^1_{ab}-\frac{1}{2})H_1} + \widetilde{u}_\alpha S^\alpha\, \sigma_{1-\theta^1_{ab}}\, e^{-i(\theta^1_{ab}+\frac{1}{2})H_1} \right)$$
$$\times\, \sigma_{1-\theta^2_{ab}}\, e^{-i(\theta^2_{ab}-\frac{1}{2})H_2}\, \sigma_{-\theta^3_{ab}}\, e^{-i(\theta^3_{ab}+\frac{1}{2})H_3}\, e^{ikX}\ , \quad (2.15)$$

where BRST invariance implies for the left- and right-moving Weyl components $\widetilde{u}_\alpha$ and $\overline{\widetilde{v}}_{\dot\alpha}$ to satisfy the Dirac equation, analogously to the Weyl components of massive fermion vertex operator (2.12).

Finally, in an analogous fashion as before for the g_ψ we determine $g_{\widetilde{\psi}}$ by computing the three-point function $\langle A\widetilde{\psi}\overline{\widetilde{\psi}}\rangle$ resulting in

$$g_{\widetilde{\psi}} = g_{\overline{\widetilde{\psi}}} = \sqrt{2}\alpha'^{\frac{3}{4}} e^{\phi_{10}/2}\ . \quad (2.16)$$

Given (2.4), (2.5), (2.12) and (2.15) we have now all the four vertex operators required to perform the computation of the disk scattering amplitude (2.3).

3. The scattering amplitude

Having derived all the vertex operators we can proceed to compute the amplitude containing two gauge bosons the two massive fermions $\widetilde{\psi}$ and $\overline{\widetilde{\psi}}$. Using the necessary correlators [7], we obtain for the scattering amplitude (2.3)

$$\mathscr{M} \sim \sqrt{\alpha'}\,\widetilde{C}_{D2}\, g_{A_x} g_{A_y}\, g^2_{\widetilde{\psi}} \int \frac{\prod_{i=1}^4 dx_i}{V_{CKG}}\, x_{12}^{s-1}\, x_{13}^{t+\theta^1_{ab}}\, x_{14}^{u+\theta^1_{ab}-1}\, x_{23}^{u+\theta^1_{ab}-1}\, x_{24}^{t+\theta^1_{ab}}\, x_{34}^{s-1}$$
$$\times \Bigg\{ \left((k_2\cdot\varepsilon_1)\varepsilon_{2\nu} - (k_1\cdot\varepsilon_2)\varepsilon_{1\nu} + (\varepsilon_1\cdot\varepsilon_2)k_{1\nu} + \frac{x_{12}\,x_{34}}{x_{13}\,x_{24}}(k_3\cdot\varepsilon_1)\varepsilon_{2\nu} \right)$$
$$\times \left((\widetilde{v}_3\,\sigma^\nu \overline{\widetilde{v}}_4) + (\overline{\widetilde{u}}_3\,\overline{\sigma}^\nu \widetilde{u}_4) \right)$$
$$+ \frac{1}{2}\frac{x_{12}\,x_{34}}{x_{13}\,x_{24}}\,\varepsilon_{1\mu}\,\varepsilon_{2\nu}\,k_{1\lambda} \left((\widetilde{v}_3\,\sigma^\lambda\overline{\sigma}^\mu\sigma^\nu\,\overline{\widetilde{v}}_4) + (\overline{\widetilde{u}}_3\,\overline{\sigma}^\lambda\sigma^\mu\overline{\sigma}^\nu\widetilde{u}_4) \right) \Bigg\}\ . \quad (3.1)$$

Here $x_{ij} = x_i - x_j$ and we dropped for the moment the Chan-Paton factors which determine the relative positions of the respective vertex operator. We should mention here that the order of the Chan Paton factors depends crucially on whether the two gauge bosons from the same, $x = y$ or from different D-brane stacks $x \neq y$.

[6] Note that here we consider the D-brane intersection *ba*, which in contrast to the system *ab* does have two negative angles in the first two complex internal dimensions and a positive intersection angle in the third complex internal dimension.

[7] For correlators involving excited bosonic twist fields, see [80–85].

Finally, the factor V_{CKG} is the volume of the conformal Killing group of the disk, which can be accounted for by fixing three vertex operator positions and adding the appropriate c-ghost correlator. While the generic four point amplitude is given by the sum over all 6 cyclic invariant orderings of the vertex operators, where each ordering gives rise to different integration region, the trace over Chan-Paton factors of the respective vertex operators gives only non-vanishing amplitudes for specific integration regions.

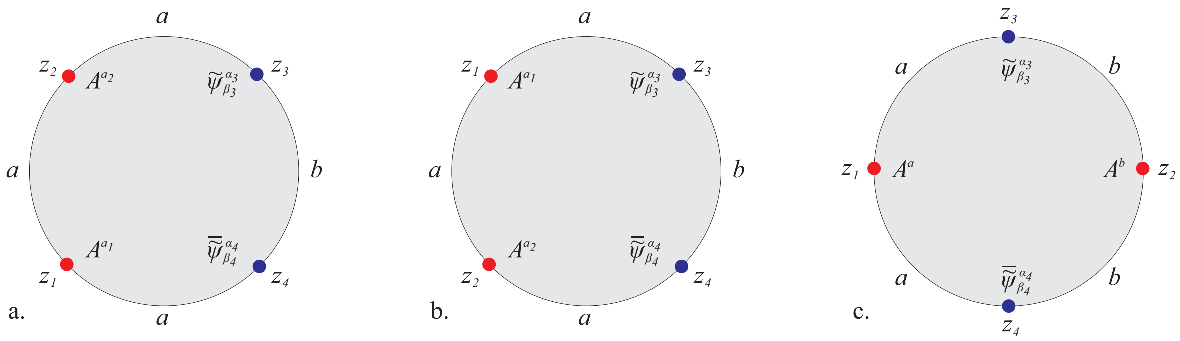

Figure 3: The possible orderings of the respective vertex operators.

Using the symmetries of the disk, we fix the vertex operator positions to

$$x_1 = 0 \qquad x_3 = 1 \qquad x_4 = \infty \tag{3.2}$$

with both gauge bosons arising from the same D-brane stack (see figure 3.a and 3.b) we have the two integration regions $-\infty < x_2 < 0$ and $0 < x_2 < 1$. On the other hand if the two gauge bosons arise from different D-brane stacks we have only one integration region, namely $1 < x_2 < \infty$ (see figure 3.c). Adding the c-ghost correlator

$$\langle c(x_1)c(x_3)c(x_4)\rangle = x_{13}x_{14}x_{34} \tag{3.3}$$

due to the fixing of the three vertex operator positions we obtain

- for the amplitude with the two gauge bosons living on the same D-brane stack

$$\begin{aligned}\mathcal{M}\left[A^{a_1}[\varepsilon_1,k_1],A^{a_2}[\varepsilon_2,k_2],\widetilde{\psi}[\widetilde{v}_3,\bar{\widetilde{v}}_3,k_3],\bar{\widetilde{\psi}}[\bar{\widetilde{v}}_4,\widetilde{u}_4,k_4]\right] &= -2\alpha' g^2_{Dp_a}\left(K+\widetilde{K}\right)\\ \times\left\{\mathrm{Tr}\left[T^{a_1}T^{a_2}T^{ab}T^{ba}\right]B[S,U]+\mathrm{Tr}\left[T^{a_2}T^{a_1}T^{ab}T^{ba}\right]\frac{T}{U}B[S,T]\right\}&\end{aligned} \tag{3.4}$$

  where the kinematic factors take the form

$$\begin{aligned}K = &\left\{(k_2\cdot\varepsilon_1)\varepsilon_{2\nu}-(k_1\cdot\varepsilon_2)\varepsilon_{1\nu}+(\varepsilon_1\cdot\varepsilon_2)k_{1\nu}-\frac{S}{T}(k_3\cdot\varepsilon_1)\varepsilon_{2\nu}\right\}(\widetilde{v}_3\,\sigma^\nu\bar{\widetilde{v}}_4)\\ &-\frac{1}{2}\frac{S}{T}\varepsilon_{1\mu}\,\varepsilon_{2\nu}\,k_{1\lambda}\left(\widetilde{v}_3\,\sigma^\lambda\bar{\sigma}^\mu\sigma^\nu\,\bar{\widetilde{v}}_4\right)\end{aligned} \tag{3.5}$$

  and $K\to\widetilde{K}$ if $(\widetilde{v}_3\,\sigma^\nu\bar{\widetilde{v}}_4)\to\left(\bar{\widetilde{u}}_3\,\bar{\sigma}^\nu\widetilde{u}_4\right)$ and $\left(\widetilde{v}_3\,\sigma^\lambda\bar{\sigma}^\mu\sigma^\nu\,\bar{\widetilde{v}}_4\right)\to\left(\bar{\widetilde{u}}_3\,\bar{\sigma}^\lambda\,\sigma^\mu\bar{\sigma}^\nu\widetilde{u}_4\right)$. Here we introduced the modified Mandelstam variables, defined as

$$S = \alpha'\,(k_1+k_2)^2 \qquad T = \alpha'\,(k_1+k_3)^2+\theta^1_{ab} \qquad U = \alpha'\,(k_1+k_4)^2+\theta^1_{ab} \tag{3.6}$$

that satisfy the relation

$$S+T+U=0\,. \tag{3.7}$$

Finally, $B[m,n]$ denotes the Euler Beta function that can be represented by the integral

$$B[m,n]=\int_0^1 x^{m-1}(1-x)^{n-1}=\frac{\Gamma[m]\,\Gamma[n]}{\Gamma[m+n]}\,. \tag{3.8}$$

- On the other hand for the two gauge bosons arising from two different stacks of D-branes we obtain

$$\begin{aligned}\mathcal{M}\left[A^a[\varepsilon_1,k_1],A^b[\varepsilon_2,k_2],\widetilde{\psi}[\widetilde{v}_3,\bar{\widetilde{v}}_3,k_3],\bar{\widetilde{\psi}}[\bar{\widetilde{v}}_4,\widetilde{u}_4,k_4]\right]&=-2\alpha'\,g_{Dp_a}\,g_{Dp_b}\left(K+\widetilde{K}\right)\\ &\times\mathrm{Tr}\left[T^aT^{ab}T^bT^{ba}\right]\frac{T}{S}B[T,U]\,.\end{aligned} \tag{3.9}$$

Note that in the limit $\theta^1_{ab}\to 0$, where the fermion $\widetilde{\psi}$ becomes massless, we get exactly the same results as for the scattering amplitude of two gauge bosons with two massless fermions as computed in [30].

After a few simple manipulations of the traces the amplitudes take the form

$$\begin{aligned}\mathcal{M}\left[A^{a_1}[\varepsilon_1,k_1],A^{a_2}[\varepsilon_2,k_2],\widetilde{\psi}[\widetilde{v}_3,\bar{\widetilde{v}}_3,k_3],\bar{\widetilde{\psi}}[\bar{\widetilde{v}}_4,\widetilde{u}_4,k_4]\right]&=2\alpha'\,g^2_{Dp_a}\left(K+\widetilde{K}\right)\\ &\times\frac{1}{U}\left\{[T^{a_1}T^{a_2}]^{\alpha_3}_{\alpha_4}\,\delta^{\beta_4}_{\beta_3}\frac{T}{S}\widehat{V}_T+[T^{a_2}T^{a_1}]^{\alpha_3}_{\alpha_4}\,\delta^{\beta_4}_{\beta_3}\frac{U}{S}\widehat{V}_U\right\}\end{aligned} \tag{3.10}$$

$$\begin{aligned}\mathcal{M}\left[A^a[\varepsilon_1,k_1],A^b[\varepsilon_2,k_2],\widetilde{\psi}[\widetilde{v}_3,\bar{\widetilde{v}}_3,k_3],\bar{\widetilde{\psi}}[\bar{\widetilde{v}}_4,\widetilde{u}_4,k_4]\right]&=2\alpha'\,g_{Dp_a}g_{Dp_b}\left(K+\widetilde{K}\right)\\ &\times\frac{1}{U}\,[T^a]^{\alpha_3}_{\alpha_4}\,[T^b]^{\beta_4}_{\beta_3}\widehat{V}_S\,,\end{aligned} \tag{3.11}$$

where we introduced the *generalized* Veneziano formfactor given by

$$\widehat{V}_S=\frac{T\,U}{T+U}B[T,U]\,,\qquad \widehat{V}_T=\frac{S\,U}{S+U}B[S,U]\,,\qquad \widehat{V}_U=\frac{S\,T}{S+T}B[S,T]\,. \tag{3.12}$$

In the following section we take the results (3.10) and (3.11) square them and sum over all polarization and gauge configurations. Eventually we apply our findings to the quark sector of a SM D-brane realization.

4. Squared amplitudes

In this chapter we provide the squared amplitudes summed over helicities and spins, as well as over the gauge indices of initial and final particles involved in the scattering process. Eventually we average over all the initial gauge as well as polarization configurations.

In order to sum over all spins s_i and helicities h_i of the massive fermions and gauge bosons, we use the completeness relations

$$\sum_{h_1,h_1'} \varepsilon_1^{*\mu} \varepsilon_1^{\nu} = -g^{\mu\nu} \tag{4.1}$$

$$\sum_{s,s'} \overline{\widetilde{v}}^{\dot{a}}(k) \widetilde{v}^{a}(k) = -k^{\mu} \overline{\sigma}_{\mu}^{\dot{a}a} \qquad \sum_{s,s'} \widetilde{u}_{a}(k) \overline{\widetilde{u}}_{\dot{a}}(k) = -k_{\mu} \sigma_{a\dot{a}}^{\mu} \tag{4.2}$$

$$\sum_{s,s'} \widetilde{v}^{a}(k) \widetilde{u}_{b}(k) = m\delta_{b}^{a} \qquad \sum_{s,s'} \overline{\widetilde{v}}^{\dot{a}}(k) \overline{\widetilde{u}}_{\dot{b}}(k) = m\delta_{\dot{b}}^{\dot{a}} \tag{4.3}$$

for the gauge bosons as well as the massive fermions one obtains applying various trace identities of sigma matrices (for more details [1]). Concerning the summation over the gauge indices we use

$$\sum_{a} T^a T^a = \frac{N^2-1}{2N} \mathbb{I}_N \, , \qquad \sum_{a_1,a_2,n} f^{a_1 a_2 n} f^{a_1 a_2 n} = N(N^2-1) \tag{4.4}$$

which holds for $SU(N)$ gauge symmetries with f^{abc} denoting the totally antisymmetric structure constants, $\left[T^a, T^b\right] = i f^{abc} T^c$.

The squared amplitudes finally become

$$\begin{aligned} \left| \mathscr{M}\left[A^a A^a \to \widetilde{\Psi}_{ab} \overline{\widetilde{\Psi}}_{ab}\right]\right|^2 &= \frac{8 g_a^4}{S^2} \mathscr{F}[S,T,U,m_{\widetilde{\Psi}_{ab}}] \\ &\times \frac{N_b\left(N_a^2-1\right)^2}{4N_a} \left\{ (T\widehat{V}_T + U\widehat{V}_U)^2 - \frac{2N_a^2}{N_a^2-1} U\widehat{V}_U T\widehat{V}_T \right\} \end{aligned} \tag{4.5}$$

$$\left| \mathscr{M}\left[A^a A^b \to \widetilde{\Psi}_{ab} \overline{\widetilde{\Psi}}_{ab}\right]\right|^2 = \frac{8 g_a^2 g_b^2}{S^2} \mathscr{F}[S,T,U,m_{\widetilde{\Psi}_{ab}}] \frac{\left(N_a^2-1\right)\left(N_b^2-1\right)}{4} \widehat{V}_S^2 \, , \tag{4.6}$$

where for simplicity we define

$$\mathscr{F}[S,T,U,m] = \frac{T}{U} + \frac{U}{T} + 4\alpha' m^2 \left(\frac{1}{T} + \frac{1}{U}\right) - 4\alpha'^2 m^4 \left(\frac{1}{T} + \frac{1}{U}\right)^2 . \tag{4.7}$$

Note in case of abelian gauge bosons the result (4.5) reproduces in the low energy limit $\alpha' \to 0$ as expected the Klein-Nishina formula, describing the Compton cross section [1].

Connection to the Standard Model

In the following we apply the formulae evaluated above to the quark and leptonic sector of a Standard Model D-brane realization. Recall that the left-handed quarks Q_L and its stringy excitations $\widetilde{Q}_L$ arise from the intersection of color D-brane stack and the $SU(2)_L$ D-brane stack. On the other hand the right-handed quarks d_R and its stringy excitations $\widetilde{d}_R$, as well as u_R and its stringy excitation $\widetilde{u}_R$ are localized at the intersection of the color D-brane stack and a $U(1)$ D-brane stack, which we call c and d, respectively [8] (see figure 4).

[8]Here we assume that the quarks are always realized as bi-fundamentals and moreover, that all three families arise from the same D-brane stack intersection. One can easily generalise this scenario to setups in which the different families are differently charged with respect to D-brane gauge symmetries. For original work on local D-brane configurations, see [86, 87]. For a systematic analysis of local D-brane configurations, see [65, 88–92].

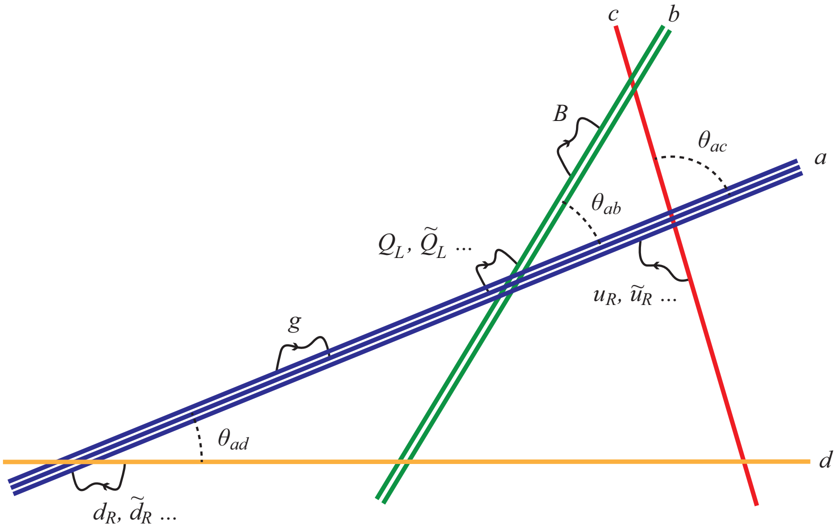

Figure 4: Quarks and their stringy excitations of a SM D-brane realization.

Using (4.5), and averaging over all initial states which implies dividing by $2(N_a^2-1)$ for each gauge boson due to the two helicities and the dimension of the adjoint representation of $SU(N_a)$, we get

$$\left|\mathscr{M}\left[gg \to \widetilde{Q}_L\overline{\widetilde{Q}}_L\right]\right|^2 = g^4\frac{1}{S^2}\mathscr{F}[S,T,U,m_{\widetilde{Q}_L}]\,\frac{1}{3}\left\{(T\widehat{V}_T+U\widehat{V}_U)^2-\frac{9}{4}U\widehat{V}_U T\widehat{V}_T\right\}\,, \tag{4.8}$$

where $\alpha' m^2_{\widetilde{Q}_L} = \theta_{ab}$ [9], g denoting the color gauge coupling, and $N_a = 3$ and $N_b = 2$ due to the fact that $\widetilde{Q}_L$ is stretched between an $SU(3)$ and $SU(2)$ D-brane stack. Similarly, we can compute scatterings like $gg \to \widetilde{d}_R\overline{\widetilde{d}}_R$ and $gg \to \widetilde{u}_R\overline{\widetilde{u}}_R$. Using also the crossing relations we can even compute similar scaterings with gluons at the final state.

Finally, we also display the processes in which a B boson, the gauge boson living on the $SU(2)$ D-brane stack, is involved. The squared amplitude can be derived from (4.6) by averaging over the $2(N_a^2-1)$ and $2(N_b^2-1)$ initial configurations of g and B, respectively. For $N_a = 3$ and $N_b = 2$ one obtains

$$\left|\mathscr{M}\left[gB \to \widetilde{Q}_L\overline{\widetilde{Q}}_L\right]\right|^2 = \frac{1}{2}g^2 g_b^2\mathscr{F}[S,T,U,m_{\widetilde{Q}_L}]\widehat{V}_S^2\,, \tag{4.9}$$

where g again denotes the color gauge coupling and g_B is the gauge coupling of the B-boson gauge group. From (4.9) we can determine all other processes using the crossing relations in an analogous fashion as done above. We display the processes in table 1 where we again average over all the initial polarization and gauge configurations.

[9] The mass of the lightest stringy excitation depends on the smallest intersection angle. In the explicit computation we assume the smallest of the three intersection angles to be θ^1_{ab}. As pointed out before the analysis is independent of the choice which of the three intersection angles is the smallest.

process	$\|\mathscr{M}\|^2$
$g\,g \to \widetilde{Q}_L\,\overline{\widetilde{Q}}_L$	$g^4\,\frac{1}{S^2}\mathscr{F}[S,T,U,m_{\widetilde{Q}_L}]\,\frac{1}{3}\left\{(T\widehat{V}_T+U\widehat{V}_U)^2-\frac{9}{4}U\widehat{V}_U T\widehat{V}_T\right\}$
$g\,g \to \widetilde{d}_R\,\overline{\widetilde{d}}_R$	$g^4\,\frac{1}{S^2}\mathscr{F}[S,T,U,m_{\widetilde{d}_R}]\,\frac{1}{6}\left\{(T\widehat{V}_T+U\widehat{V}_U)^2-\frac{9}{4}U\widehat{V}_U T\widehat{V}_T\right\}$
$g\,g \to \widetilde{u}_R\,\overline{\widetilde{u}}_R$	$g^4\,\frac{1}{S^2}\mathscr{F}[S,T,U,m_{\widetilde{u}_R}]\,\frac{1}{6}\left\{(T\widehat{V}_T+U\widehat{V}_U)^2-\frac{9}{4}U\widehat{V}_U T\widehat{V}_T\right\}$
$g\,\widetilde{Q}_L \to g\,\overline{\widetilde{Q}}_L$	$g^4\,\frac{1}{T^2}\mathscr{F}[T,S,U,m_{\widetilde{Q}_L}]\,\frac{4}{9}\left\{(S\widehat{V}_S+U\widehat{V}_U)^2-\frac{9}{4}U\widehat{V}_U S\widehat{V}_S\right\}$
$g\,\widetilde{d}_R \to g\,\overline{\widetilde{d}}_R$	$g^4\,\frac{1}{T^2}\mathscr{F}[T,S,U,m_{\widetilde{d}_R}]\,\frac{4}{9}\left\{(S\widehat{V}_S+U\widehat{V}_U)^2-\frac{9}{4}U\widehat{V}_U S\widehat{V}_S\right\}$
$g\,\widetilde{u}_R \to g\,\overline{\widetilde{u}}_R$	$g^4\,\frac{1}{T^2}\mathscr{F}[T,S,U,m_{\widetilde{u}_R}]\,\frac{4}{9}\left\{(S\widehat{V}_S+U\widehat{V}_U)^2-\frac{9}{4}U\widehat{V}_U S\widehat{V}_S\right\}$
$g\,\widetilde{Q}_L \to B\,\overline{\widetilde{Q}}_L$	$g^2 g_B^2\mathscr{F}[T,S,U,m_{\widetilde{Q}_L}]\,\frac{1}{4}\widehat{V}_T^2$
$\widetilde{Q}_L\,\overline{\widetilde{Q}}_L \to g\,g$	$g^4\,\frac{1}{S^2}\mathscr{F}[S,T,U,m_{\widetilde{Q}_L}]\,\frac{16}{27}\left\{(T\widehat{V}_T+U\widehat{V}_U)^2-\frac{9}{4}U\widehat{V}_U T\widehat{V}_T\right\}$
$\widetilde{d}_R\,\overline{\widetilde{d}}_R \to g\,g$	$g^4\,\frac{1}{S^2}\mathscr{F}[S,T,U,m_{\widetilde{d}_R}]\,\frac{32}{27}\left\{(T\widehat{V}_T+U\widehat{V}_U)^2-\frac{9}{4}U\widehat{V}_U T\widehat{V}_T\right\}$
$\widetilde{u}_R\,\overline{\widetilde{u}}_R \to g\,g$	$g^4\,\frac{1}{S^2}\mathscr{F}[S,T,U,m_{\widetilde{u}_R}]\,\frac{32}{27}\left\{(T\widehat{V}_T+U\widehat{V}_U)^2-\frac{9}{4}U\widehat{V}_U T\widehat{V}_T\right\}$
$\widetilde{Q}_L\,\overline{\widetilde{Q}}_L \to g\,B$	$g^2 g_B^2\mathscr{F}[S,T,U,m_{\widetilde{Q}_L}]\,\frac{1}{3}\widehat{V}_S^2$

Table 1: Processes involving two gauge bosons and stringy excitations of quarks.

5. Conclusions

In [1] we study stringy excitations of states localized at the intersection of two D-brane stacks. We determine the vertex operators of these states and we compute the disk diagram containing two gauge bosons and two light stringy states. We sum over all polarizations and color configurations of the initial and final states. In the limit of $\theta_{ab} \to 0$, i.e. in the limit in which those light stringy states become massless, we recover the results of [30], which analyzed the scattering of two gauge bosons onto two massless fermions. In addition in case the gauge bosons are $U(1)$ gauge bosons, we show that the low energy limit our squared amplitude reproduces with the Klein-Nishina formula. Finally, we apply our results to various scatterings of two SM gauge bosons onto two stringy excitations of the quark sector.

Extensive study of decays of such light stringy excitations of the SM matter sector is very interesting since such light stringy states can be easily distinguished from other KK scenarios and might be visible at LHC.

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References

[1] P. Anastasopoulos and R. Richter, *Production of light stringy states*, *JHEP* **1412** (2014) 059, [arXiv:1408.4810].

[2] R. Blumenhagen, M. Cvetič, P. Langacker, and G. Shiu, *Toward realistic intersecting D-brane models*, *Ann.Rev.Nucl.Part.Sci.* **55** (2005) 71–139, [hep-th/0502005].

[3] R. Blumenhagen, B. Körs, D. Lüst, and S. Stieberger, *Four-dimensional String Compactifications with D-Branes, Orientifolds and Fluxes*, *Phys.Rept.* **445** (2007) 1–193, [hep-th/0610327].

[4] F. Marchesano, *Progress in D-brane model building*, *Fortsch.Phys.* **55** (2007) 491–518, [hep-th/0702094].

[5] M. Cvetič and J. Halverson, *TASI Lectures: Particle Physics from Perturbative and Non-perturbative Effects in D-braneworlds*, arXiv:1101.2907.

[6] N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, *The Hierarchy problem and new dimensions at a millimeter*, *Phys.Lett.* **B429** (1998) 263–272, [hep-ph/9803315].

[7] I. Antoniadis, S. Dimopoulos, and G. Dvali, *Millimeter range forces in superstring theories with weak scale compactification*, *Nucl.Phys.* **B516** (1998) 70–82, [hep-ph/9710204].

[8] I. Antoniadis, N. Arkani-Hamed, S. Dimopoulos, and G. Dvali, *New dimensions at a millimeter to a Fermi and superstrings at a TeV*, *Phys.Lett.* **B436** (1998) 257–263, [hep-ph/9804398].

[9] **CMS Collaboration** Collaboration, S. Chatrchyan *et. al.*, *Search for Resonances in the Dijet Mass Spectrum from 7 TeV pp Collisions at CMS*, *Phys.Lett.* **B704** (2011) 123–142, [arXiv:1107.4771].

[10] **ATLAS Collaboration** Collaboration, G. Aad *et. al.*, *ATLAS search for new phenomena in dijet mass and angular distributions using pp collisions at $\sqrt{s} = 7$ TeV*, *JHEP* **1301** (2013) 029, [arXiv:1210.1718].

[11] **CMS Collaboration** Collaboration, S. Chatrchyan *et. al.*, *Search for narrow resonances using the dijet mass spectrum in pp collisions at $\sqrt{s}$=8??TeV*, *Phys.Rev.* **D87** (2013), no. 11 114015, [arXiv:1302.4794].

[12] E. Dudas and J. Mourad, *String theory predictions for future accelerators*, *Nucl.Phys.* **B575** (2000) 3–34, [hep-th/9911019].

[13] E. Accomando, I. Antoniadis, and K. Benakli, *Looking for TeV scale strings and extra dimensions*, *Nucl.Phys.* **B579** (2000) 3–16, [hep-ph/9912287].

[14] S. Cullen, M. Perelstein, and M. E. Peskin, *TeV strings and collider probes of large extra dimensions*, *Phys.Rev.* **D62** (2000) 055012, [hep-ph/0001166].

[15] E. Kiritsis and P. Anastasopoulos, *The Anomalous magnetic moment of the muon in the D-brane realization of the standard model*, *JHEP* **0205** (2002) 054, [hep-ph/0201295].

[16] I. Antoniadis, E. Kiritsis, and J. Rizos, *Anomalous U(1)s in type 1 superstring vacua*, *Nucl.Phys.* **B637** (2002) 92–118, [hep-th/0204153].

[17] D. Ghilencea, L. Ibáñez, N. Irges, and F. Quevedo, *TeV scale Z-prime bosons from D-branes*, *JHEP* **0208** (2002) 016, [hep-ph/0205083].

[18] P. Anastasopoulos, *4-D anomalous U(1)'s, their masses and their relation to 6-D anomalies*, *JHEP* **0308** (2003) 005, [hep-th/0306042].

[19] P. Anastasopoulos, *Anomalous U(1)s masses in nonsupersymmetric open string vacua*, *Phys.Lett.* **B588** (2004) 119–126, [hep-th/0402105].

[20] P. Anastasopoulos, *Orientifolds, anomalies and the standard model*, hep-th/0503055.

[21] C. Burgess, J. Matias, and F. Quevedo, *MSLED: A Minimal supersymmetric large extra dimensions scenario*, *Nucl.Phys.* **B706** (2005) 71–99, [hep-ph/0404135].

[22] P. Burikham, T. Figy, and T. Han, *TeV-scale string resonances at hadron colliders*, *Phys.Rev.* **D71** (2005) 016005, [hep-ph/0411094].

[23] D. Chialva, R. Iengo, and J. G. Russo, *Cross sections for production of closed superstrings at high energy colliders in brane world models*, *Phys.Rev.* **D71** (2005) 106009, [hep-ph/0503125].

[24] C. Corianò, N. Irges, and E. Kiritsis, *On the effective theory of low scale orientifold string vacua*, *Nucl.Phys.* **B746** (2006) 77–135, [hep-ph/0510332].

[25] M. Bianchi and A. V. Santini, *String predictions for near future colliders from one-loop scattering amplitudes around D-brane worlds*, *JHEP* **0612** (2006) 010, [hep-th/0607224].

[26] P. Anastasopoulos, M. Bianchi, E. Dudas, and E. Kiritsis, *Anomalies, anomalous U(1)'s and generalized Chern-Simons terms*, *JHEP* **0611** (2006) 057, [hep-th/0605225].

[27] L. A. Anchordoqui, H. Goldberg, S. Nawata, and T. R. Taylor, *Jet signals for low mass strings at the LHC*, *Phys.Rev.Lett.* **100** (2008) 171603, [arXiv:0712.0386].

[28] P. Anastasopoulos, F. Fucito, A. Lionetto, G. Pradisi, A. Racioppi, and Y. Stanev, *Minimal Anomalous U(1)-prime Extension of the MSSM*, *Phys.Rev.* **D78** (2008) 085014, [arXiv:0804.1156].

[29] L. A. Anchordoqui, H. Goldberg, S. Nawata, and T. R. Taylor, *Direct photons as probes of low mass strings at the CERN LHC*, *Phys.Rev.* **D78** (2008) 016005, [arXiv:0804.2013].

[30] D. Lüst, S. Stieberger, and T. R. Taylor, *The LHC String Hunter's Companion*, *Nucl.Phys.* **B808** (2009) 1–52, [arXiv:0807.3333].

[31] L. A. Anchordoqui, H. Goldberg, D. Lüst, S. Nawata, S. Stieberger, and T. Taylor, *Dijet signals for low mass strings at the LHC*, *Phys.Rev.Lett.* **101** (2008) 241803, [arXiv:0808.0497].

[32] R. Armillis, C. Coriano, M. Guzzi, and S. Morelli, *An Anomalous Extra Z Prime from Intersecting Branes with Drell-Yan and Direct Photons at the LHC*, *Nucl.Phys.* **B814** (2009) 156–179, [arXiv:0809.3772].

[33] F. Fucito, A. Lionetto, A. Mammarella, and A. Racioppi, *Stueckelino dark matter in anomalous U(1)-prime models*, *Eur.Phys.J.* **C69** (2010) 455–465, [arXiv:0811.1953].

[34] L. A. Anchordoqui, H. Goldberg, D. Lüst, S. Nawata, S. Stieberger, and T. Taylor, *LHC Phenomenology for String Hunters*, *Nucl.Phys.* **B821** (2009) 181–196, [arXiv:0904.3547].

[35] L. A. Anchordoqui, H. Goldberg, D. Lüst, S. Stieberger, and T. R. Taylor, *String Phenomenology at the LHC*, *Mod.Phys.Lett.* **A24** (2009) 2481–2490, [arXiv:0909.2216].

[36] D. Lüst, O. Schlotterer, S. Stieberger, and T. Taylor, *The LHC String Hunter's Companion (II): Five-Particle Amplitudes and Universal Properties*, *Nucl.Phys.* **B828** (2010) 139–200, [arXiv:0908.0409].

[37] Z. Dong, T. Han, M.-x. Huang, and G. Shiu, *Top Quarks as a Window to String Resonances*, *JHEP* **1009** (2010) 048, [arXiv:1004.5441].

[38] W.-Z. Feng, D. Lüst, O. Schlotterer, S. Stieberger, and T. R. Taylor, *Direct Production of Lightest Regge Resonances*, *Nucl.Phys.* **B843** (2011) 570–601, [arXiv:1007.5254].

[39] L. A. Anchordoqui, W.-Z. Feng, H. Goldberg, X. Huang, and T. R. Taylor, *Searching for string resonances in $e^+ e^-$ and $\gamma\gamma$ collisions*, *Phys.Rev.* **D83** (2011) 106006, [arXiv:1012.3466].

[40] M. Cicoli, C. Burgess, and F. Quevedo, *Anisotropic Modulus Stabilisation: Strings at LHC Scales with Micron-sized Extra Dimensions*, *JHEP* **1110** (2011) 119, [arXiv:1105.2107].

[41] D. Carmi, *TeV Scale Strings and Scattering Amplitudes at the LHC*, arXiv:1109.5161.

[42] L. A. Anchordoqui, I. Antoniadis, H. Goldberg, X. Huang, D. Lüst, T. Taylor, and B. Vlcek, *LHC Phenomenology and Cosmology of String-Inspired Intersecting D-Brane Models*, *Phys.Rev.* **D86** (2012) 066004, [arXiv:1206.2537].

[43] L. A. Anchordoqui, I. Antoniadis, H. Goldberg, X. Huang, D. Lüst, and T. Taylor, *Z'-gauge Bosons as Harbingers of Low Mass Strings*, *Phys.Rev.* **D85** (2012) 086003, [arXiv:1107.4309].

[44] L. A. Anchordoqui, H. Goldberg, X. Huang, D. Lüst, and T. R. Taylor, *Stringy origin of Tevatron Wjj anomaly*, *Phys.Lett.* **B701** (2011) 224–228, [arXiv:1104.2302].

[45] M. Hashi and N. Kitazawa, *Detectability of the second resonance of low-scale string models at the LHC*, *JHEP* **1303** (2013) 127, [arXiv:1212.5372].

[46] D. Chialva, P. B. Dev, and A. Mazumdar, *Multiple dark matter scenarios from ubiquitous stringy throats*, *Phys.Rev.* **D87** (2013), no. 6 063522, [arXiv:1211.0250].

[47] D. Lüst and T. R. Taylor, *Limits on Stringy Signals at the LHC*, arXiv:1308.1619.

[48] L. A. Anchordoqui, I. Antoniadis, D.-C. Dai, W.-Z. Feng, H. Goldberg, *et. al.*, *String Resonances at Hadron Colliders*, arXiv:1407.8120.

[49] D. Berenstein, *TeV-Scale strings*, arXiv:1401.4491.

[50] P. Anastasopoulos, M. Bianchi, and R. Richter, *Light stringy states*, *JHEP* **1203** (2012) 068, [arXiv:1110.5424].

[51] R. Blumenhagen, L. Görlich, B. Körs, and D. Lüst, *Noncommutative compactifications of type I strings on tori with magnetic background flux*, *JHEP* **0010** (2000) 006, [hep-th/0007024].

[52] C. Angelantonj, I. Antoniadis, E. Dudas, and A. Sagnotti, *Type I strings on magnetized orbifolds and brane transmutation*, *Phys.Lett.* **B489** (2000) 223–232, [hep-th/0007090].

[53] G. Aldazabal, S. Franco, L. E. Ibáñez, R. Rabadán, and A. Uranga, *Intersecting brane worlds*, *JHEP* **0102** (2001) 047, [hep-ph/0011132].

[54] G. Aldazabal, S. Franco, L. E. Ibáñez, R. Rabadán, and A. Uranga, *D = 4 chiral string compactifications from intersecting branes*, *J.Math.Phys.* **42** (2001) 3103–3126, [hep-th/0011073].

[55] S. Förste, G. Honecker, and R. Schreyer, *Supersymmetric Z(N) x Z(M) orientifolds in 4-D with D branes at angles*, *Nucl.Phys.* **B593** (2001) 127–154, [hep-th/0008250].

[56] L. E. Ibáñez, F. Marchesano, and R. Rabádan, *Getting just the standard model at intersecting branes*, *JHEP* **0111** (2001) 002, [hep-th/0105155].

[57] M. Cvetič, G. Shiu, and A. M. Uranga, *Three family supersymmetric standard - like models from intersecting brane worlds*, *Phys.Rev.Lett.* **87** (2001) 201801, [hep-th/0107143].

[58] M. Cvetič, G. Shiu, and A. M. Uranga, *Chiral four-dimensional N=1 supersymmetric type 2A orientifolds from intersecting D6 branes*, *Nucl.Phys.* **B615** (2001) 3–32, [hep-th/0107166].

[59] G. Honecker, *Chiral supersymmetric models on an orientifold of Z(4) x Z(2) with intersecting D6-branes*, *Nucl.Phys.* **B666** (2003) 175–196, [hep-th/0303015].

[60] M. Cvetič, P. Langacker, T.-j. Li, and T. Liu, *D6-brane splitting on type IIA orientifolds*, *Nucl.Phys.* **B709** (2005) 241–266, [hep-th/0407178].

[61] G. Honecker, *Chiral N=1 4-D orientifolds with D-branes at angles*, *Mod.Phys.Lett.* **A19** (2004) 1863–1879, [hep-th/0407181].

[62] G. Honecker and T. Ott, *Getting just the supersymmetric standard model at intersecting branes on the Z(6) orientifold*, *Phys.Rev.* **D70** (2004) 126010, [hep-th/0404055].

[63] R. Blumenhagen, F. Gmeiner, G. Honecker, D. Lüst, and T. Weigand, *The Statistics of supersymmetric D-brane models*, *Nucl.Phys.* **B713** (2005) 83–135, [hep-th/0411173].

[64] R. Blumenhagen, M. Cvetič, F. Marchesano, and G. Shiu, *Chiral D-brane models with frozen open string moduli*, *JHEP* **0503** (2005) 050, [hep-th/0502095].

[65] F. Gmeiner, R. Blumenhagen, G. Honecker, D. Lüst, and T. Weigand, *One in a billion: MSSM-like D-brane statistics*, *JHEP* **0601** (2006) 004, [hep-th/0510170].

[66] D. Bailin and A. Love, *Towards the supersymmetric standard model from intersecting D6-branes on the Z-prime(6) orientifold*, *Nucl.Phys.* **B755** (2006) 79–111, [hep-th/0603172].

[67] M. Cvetič and P. Langacker, *New Grand Unified Models with Intersecting D6-branes, Neutrino Masses, and Flipped SU(5)*, *Nucl.Phys.* **B776** (2007) 118–137, [hep-th/0607238].

[68] C.-M. Chen, T. Li, V. Mayes, and D. V. Nanopoulos, *A Realistic world from intersecting D6-branes*, *Phys.Lett.* **B665** (2008) 267–270, [hep-th/0703280].

[69] D. Bailin and A. Love, *Almost the supersymmetric standard model from intersecting D6-branes on the Z(6)-prime orientifold*, *Phys.Lett.* **B651** (2007) 324–328, [arXiv:0705.0646].

[70] F. Gmeiner and G. Honecker, *Mapping an Island in the Landscape*, *JHEP* **0709** (2007) 128, [arXiv:0708.2285].

[71] D. Bailin and A. Love, *Constructing the supersymmetric Standard Model from intersecting D6-branes on the Z(6)-prime orientifold*, *Nucl.Phys.* **B809** (2009) 64–109, [arXiv:0801.3385].

[72] F. Gmeiner and G. Honecker, *Millions of Standard Models on Z-prime(6)?*, *JHEP* **0807** (2008) 052, [arXiv:0806.3039].

[73] G. Honecker, M. Ripka, and W. Staessens, *The Importance of Being Rigid: D6-Brane Model Building on $T^6/Z_2xZ'_6$ with Discrete Torsion*, *Nucl.Phys.* **B868** (2013) 156–222, [arXiv:1209.3010].

[74] J. Polchinski, *String theory. Vol. 1: An introduction to the bosonic string*, .

[75] J. Polchinski, *String theory. Vol. 2: Superstring theory and beyond*, .

[76] M. Cvetic and R. Richter, *Proton decay via dimension-six operators in intersecting D6-brane models*, *Nucl.Phys.* **B762** (2007) 112–147, [hep-th/0606001].

[77] M. Bertolini, M. Billó, A. Lerda, J. F. Morales, and R. Russo, *Brane world effective actions for D-branes with fluxes*, *Nucl.Phys.* **B743** (2006) 1–40, [hep-th/0512067].

[78] M. Cvetic, I. Garcia-Etxebarria, and R. Richter, *Branes and instantons intersecting at angles*, *JHEP* **1001** (2010) 005, [arXiv:0905.1694].

[79] J. R. David, *Tachyon condensation using the disc partition function*, *JHEP* **07** (2001) 009, [hep-th/0012089].

[80] D. Lüst, P. Mayr, R. Richter, and S. Stieberger, *Scattering of gauge, matter, and moduli fields from intersecting branes*, *Nucl.Phys.* **B696** (2004) 205–250, [hep-th/0404134].

[81] I. Pesando, *The generating function of amplitudes with N twisted and M untwisted states*, arXiv:1107.5525.

[82] I. Pesando, *Green functions and twist correlators for N branes at angles*, *Nucl.Phys.* **B866** (2013) 87–123, [arXiv:1206.1431].

[83] P. Anastasopoulos, M. Bianchi, and R. Richter, *On closed-string twist-field correlators and their open-string descendants*, arXiv:1110.5359.

[84] P. Anastasopoulos, M. D. Goodsell, and R. Richter, *Three- and Four-point correlators of excited bosonic twist fields*, *JHEP* **1310** (2013) 182, [arXiv:1305.7166].

[85] I. Antoniadis, E. Kiritsis, and T. Tomaras, *A D-brane alternative to unification*, *Phys.Lett.* **B486** (2000) 186–193, [hep-ph/0004214].

[86] G. Aldazabal, L. E. Ibanez, F. Quevedo, and A. Uranga, *D-branes at singularities: A Bottom up approach to the string embedding of the standard model*, *JHEP* **0008** (2000) 002, [hep-th/0005067].

[87] P. Anastasopoulos, T. Dijkstra, E. Kiritsis, and A. Schellekens, *Orientifolds, hypercharge embeddings and the Standard Model*, *Nucl.Phys.* **B759** (2006) 83–146, [hep-th/0605226].

[88] M. Cvetic, J. Halverson, and R. Richter, *Realistic Yukawa structures from orientifold compactifications*, *JHEP* **0912** (2009) 063, [arXiv:0905.3379].

[89] M. Cvetic, J. Halverson, and R. Richter, *Mass Hierarchies from MSSM Orientifold Compactifications*, *JHEP* **1007** (2010) 005, [arXiv:0909.4292].

[90] M. Cvetic, J. Halverson, P. Langacker, and R. Richter, *The Weinberg Operator and a Lower String Scale in Orientifold Compactifications*, *JHEP* **1010** (2010) 094, [arXiv:1001.3148].

[91] P. Anastasopoulos, G. Leontaris, R. Richter, and A. Schellekens, *SU(5) D-brane realizations, Yukawa couplings and proton stability*, *JHEP* **1012** (2010) 011, [arXiv:1010.5188].

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