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String theory and fundamental forces/Théorie des cordes et forces fondamentales

Effective actions of intersecting brane worlds, fluxes and gaugino condensation

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Abstract

In this article we will first discuss the construction of brane world models being built either by intersecting D6-branes in type IIA orientifolds or, in the T-dual mirror picture, by D3- plus D7-branes with f-flux in type IIB orientifolds. We will show how their effective action is obtained by the calculation of scattering amplitudes between open and closed string states on intersecting D6-branes respectively on D3- and D7-branes. Secondly, turning on type IIB 3-form fluxes we will compute the induced soft supersymmetry breaking terms for the matter fields, like gaugino and scalar field masses. Finally, we will discuss the generation of 3-form flux in type IIB supergravity, which can be associated to the dynamical formation of a gaugino condensate in the confining phase of the dual $N = 1^*$ gauge theory. ***To cite this article: D. Lüst, C. R. Physique 5 (2004).***

Résumé

Action effective pour l'intersection de mondes branaires, flux et condensation de jaugino. Dans cet article nous discutons tout d'abord la construction de modèles de monde branaires construits soit par intersection de branes D6 dans des orientifolds de type IIA ou, dans la représentation T-duale, par des branes de type D3 et D7 avec des flux f dans les orientifolds de type IIB. Nous montrons comment obtenir leurs actions effectives en calculant les amplitudes de diffusion sur des intersections de branes de D6 et aussi sur des branes de type D3 et D7. Ensuite, nous allumons des flux pour la 3-forme de type IIB et nous calculons les termes de brisure douce de la supersymétrie pour les champs de matière, comme les masses du jaugino et des champs scalaires. Enfin, nous discutons la génération de flux pour la 3-forme de type IIB en supergravité, qui peut-être associée à la dynamique de la formation de condensat de jaugino dans la phase confinante de la théorie de jauge duale $\mathcal{N} = 1^*$. ***Pour citer cet article : D. Lüst, C. R. Physique 5 (2004).***

Keywords: Brane world; Gaugino; Gauge theory

Mots-clés : Monde branaires ; Jaugino ; Théorie de jauge

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1. Introduction

One of the main goals of superstring theory is the embedding of the Standard Model into a unified description of gravitational and gauge forces. In this article we will consider $N = 1$ orientifold (type I) models with D-branes (for a review see [1]). One of their characteristic features is that gravity is mediated in the entire 10-dimensional bulk by the exchange of closed strings. On the other hand, the gauge and matter fields are localized at the D-brane (intersections) and correspond to open string excitations. Of course, at the end of all realistic model building efforts $N = 1$ supersymmetry must be broken. There are (at least) two scenarios of how supersymmetry breaking can be realized: (i) The closed string sector preserves $N = 1$ SUSY, however the various open string sectors break SUSY. In order to solve the hierarchy problem, this scenario normally requires the existence of large extra dimensions, transversal to the D-branes of the Standard Model. (ii) In the second avenue, which we like to follow here, all open string sectors preserve $N = 1$ SUSY, i.e. all D-branes together with the orientifold planes, are mutually supersymmetric (either $1/2$ oder $1/4$ BPS configurations). Then the closed string sector breaks SUSY, which manifests itself as soft SUSY breaking terms in the effective action of the open string matter fields. The concrete scenario, which we will discuss in the following, is that SUSY is 'spontaneously' broken in the closed string sector by internal background fluxes of closed string field strength fields $\langle G_{ijk}\rangle \neq 0$ [2].

The further plan of the article is the following: In the next section we will introduce the type IIA intersecting D6-brane orientifold models and their IIB mirrors, which contain D9/D5 and D3/D7, respectively, with open string 2-form f-flux on the worldvolume of the D-branes [3–9]. After that, the scattering amplitudes of gauge, matter (ϕ) and moduli fields (M) from (intersecting) D-branes are computed [10]. They give rise to the $N = 1$ supersymmetric low-energy supergravity action of the matter fields, described by the Kähler potential $K(\phi, \bar{\phi}, M, \bar{M})$ and the gauge kinetic function $f(M)$. Alternatively, the effective action of orientifolds with D3/D7-branes can be obtained by dimensional reduction of the Born–Infeld action [11,12]. In the fourth section we will add internal 3-form fluxes in orientifolds with D3/D7-branes [13–21]. They will give to an effective superpotential which generically leads to spontaneous supersymmetry breaking (F-term breaking). If G_3 is a ISD $(0, 3)$-form, this case corresponds to a non-vanishing auxiliary field F^T of the overall Kähler modulus of the internal space, and supersymmetry is spontaneously broken; if G_3 is an imaginary anti-self dual (IASD) $(3, 0)$-form it is equivalent to an auxiliary field F^S. As an effect of the supersymmetry breaking in the bulk sector by F^T and/or F^S, soft supersymmetry breaking terms for the open string matter fields are induced [22–26]. For D7-branes with non-vanishing f-flux, i.e. with mixed D/N boundary conditions, the gauge kinetic function contains both the dilaton S as well as the Kähler moduli T^i. It follows that the corresponding gaugino masses get contributions both from the $(0, 3)$ G-flux, which corresponds to $F^T \neq 0$, and also from the $(3, 0)$-flux, i.e. $F^S \neq 0$, as it will happen in realistic models with 3 chiral generations. For the scalars living on the D3-branes, a mass is only generated by the $(3, 0)$-flux, whereas scalars on pure D7-branes get their masses entirely from (0,3)-flux. On the other hand, scalars on D7-branes with f-fluxes get mass contributions both from $(3, 0)$- and (0,3)-fluxes. Most importantly, 'chiral' scalar fields, which correspond to twisted open string sectors, i.e. open strings which stretch either between D3-branes and D7-branes with f-flux or two D7-branes with different type of f-flux boundary conditions, get also soft masses from $(3, 0)$- as well as from (0,3)-fluxes. In the last section we [27] will discuss the relation between $(3, 0)$ G-flux and the dynamical formation of gaugino condensation in the $N = 1^*$ model of Polchinski/Strassler [28] using the AdS/CFT correspondence. Specifically we will extend their type IIB supergravity solution with 3-form flux to third order in the mass perturbation parameter, and show that this 3-form flux corresponds in the dual field theory to a gaugino condensate.

2. IIA intersecting D6-brane models and their IIB mirrors with D3- and D7-branes

Intersecting D6-branes in type IIA orientifold compactifications can give rise to 4-dimensional models with spectra very close to the Standard Model. Let us start to consider a local D6-brane configuration in flat 10-dimensional Minkowski space-time $\mathbb{R}^{1,9}$ [6,7,29,30], see Fig. 1. Stack a, the so called color branes, consists of 3 D6-branes, the stack b, the weak branes, contains 2 D6-branes, and furthermore there are two additional c and d D6-branes. The corresponding gauge group is $U(3) \times U(2) \times U(1)^2$. The weak hypercharge group $U(1)_Y$ is a suitable linear combination of the four $U(1)$s. Note that in a compact model (see below) some of the $U(1)$ gauge bosons may get a mass by a generalized Green–Schwarz mechanism. The chiral matter fields with SM gauge quantum numbers are localized at the D-brane intersections which all fill a common $(1 + 3)$-dimensional subspace of $\mathbb{R}^{1,9}$ and are point-like in the remaining transversal 6 spational dimensions. $N = 1$ space-time supersymmetry is preserved if all stacks of D-branes preserve the following angle conditions among each other,

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

where each of the 3 intersection angles θ^i_{ab} defines the D6-brane intersection in a two-dimensional subspace $\mathbb{R}^2_i$ of the transversal 6-dimensional space.

Table 1
Topological intersection numbers

Sector	Rep.	Number
$a'a$	A_a	$\frac{1}{2}(\pi'_a \circ \pi_a + \pi_{O6} \circ \pi_a)$
$a'a$	S_a	$\frac{1}{2}(\pi'_a \circ \pi_a - \pi_{O6} \circ \pi_a)$
ab	$(\bar{N}_a, N_b)$	$\pi_a \circ \pi_b$
$a'b$	(N_a, N_b)	$\pi'_a \circ \pi_b$

Next we have to embed the D6-branes into a compact 6-dimensional space M^6, namely more precisely we are considering type IIA orientifold compactifications on a space-time

$$(\mathbb{R}^{1,3} \times M^6)/(\Omega \times I_6),$$

where Ω is the world-sheet parity group, and I_6 is a reflection on 3 internal coordinates. The fixed locus of I_6 is the orientifold 6-plane, which is supersymmetric 3-cycle π_{O6} inside M^6.

The D6$_a$-branes are wrapped around supersymmetric 3-cycles π_a, π'_a inside M^6, i.e. the D6-brane world volume has the form $\mathbb{R}^{1,3} \times \pi_a$ or $\mathbb{R}^{1,3} \times \pi'_a$, where π'_a is the I_6 reflected 3-cycle. Since we are considering D-branes and orientifold planes on a compact space, there are the following important differences compared to the flat, non-compact case:

– The 3-cycles can intersect more than once in M^6. The corresponding intersection number has a natural geometric interpretation of a family number:

$$N_F = I_{ab} \equiv \#(\pi_a \cap \pi_b) \equiv \pi_a \circ \pi_b.$$

– The internal Ramond charges of the D-branes and the orientifold plane on the compact space must cancel (Gauss law). This condition can be phrased as an equation for the homology cycles

$$\sum_a N_a(\pi_a + \pi'_a) - 4\pi_{O6} = 0. \tag{2}$$

– In order that the D-branes are stable, all NS tadpoles must cancel, which is equivalent to the cancellation of internal brane tensions. This condition can be formulated using a D-term scalar potential

$$V_D(S, U_i) = \frac{\Im S}{\sqrt{\mathrm{Vol}(M^6)}}\left(\sum_{a,a'} N_a \ \mathrm{Vol}_{\pi_a}(U_i) - 4\mathrm{Vol}_{\pi_{O6}}(U_i)\right). \tag{3}$$

Absence of the NS-tadpoles means: $\partial V_D/\partial(\Im S) = \partial V_D/\partial U_i = 0$. This is automatically ensured if all D6-branes mutually satisfy the $N = 1$ angle condition. Note the minimization of V_D fixes some of the IIA bulk complex structure moduli U_i of M^6 as well as the 4-dimensional dilaton, $\Im S = e^{-\phi_4}$.

The chiral matter spectrum is completely fixed by the topological intersection numbers of the 3-cycles of the configuration [9], as given in Table 1.

The non-Abelian gauge anomalies will cancel after satisfying the tadpole conditions and $U(1)$ anomalies are canceled by a generalized Green–Schwarz mechanism involving dimensionally reduced RR-forms.

Let us now switch to the T-dual type IIB mirror picture. The internal compact space is described by the mirror manifold $\tilde{M}^6$, and the orientifold is defined in terms of reflection group I_n, with $n = 0, 2, 4, 6$ depending on the action of the mirror symmetry. Starting from supersymmetric intersecting D6-branes, there are two possible type IIB mirror orientifolds with the following brane content:

– D5- and D9$_a$-branes, plus O9/O5-planes ($n = 0, 4$);
– D7$_a$- and D3-branes, plus O3/O7-planes ($n = 6, 2$).

The intersection angle θ_{ab} between two D6-branes is mapped to open string 2-form flux F_{ab} through 2-cycles on the D9$_a$ respectively D7$_a$-brane world volumes. Now chiral matter originates from open strings between D5–D9$_a$, D9$_a$–D9$_b$ respectively D3–D7$_a$, D7$_a$–D7$_b$ with 2-form flux. Note that V_D now fixes the IIB bulk Kähler moduli T_i of the mirror manifold $\tilde{M}^6$.

As an example of a class of type IIB orientifolds let us consider orbifold compactifications with

$$M^6 = (T^2)^3/(\mathbb{Z}_N \times \mathbb{Z}_M). \tag{4}$$

There will be always 3 Kähler moduli T^i for each of the 3 subtori T_i^2, and, depending on the orbifold group, also some unfixed complex structure moduli U^i. In addition, there is the complex dilaton field S. Next we will include D7-branes together with their open string sectors. To obtain a chiral spectrum, we must introduce (magnetic) two-form fluxes $F^j dx^j \wedge dy^j$ on the internal part of the D7-brane world volume. Together with the internal NS B-field b^j we form the complete 2-form flux

$$F = \sum_{j=1}^{3} F^j := \sum_{j=1}^{3} \left(b^j + 2\pi f^j\right) dx^j \wedge dy^j.$$

The latter gives rise to the total internal antisymmetric background

$$\begin{pmatrix} 0 & f^j \\ -f^j & 0 \end{pmatrix}, \quad f^j = \frac{1}{(2\pi)^2} \int\limits_{T^{2,j}} F^j, \tag{5}$$

w.r.t. the jth internal plane. The 2-form fluxes F^j have to obey the quantization rule:

$$m^j \frac{1}{(2\pi)^2 \alpha'} \int\limits_{T^{2,j}} F^j = n^j, \quad n \in \mathbb{Z}, \tag{6}$$

i.e. $f^j = \alpha' \frac{n^j}{m^j}$. This setup is T-dual to intersecting D6-branes in type IIA orientifold compactifications. In a compact model, all tadpoles arising from the Ramond forms C_4 and C_8 must be cancelled either by the D-branes or by the 3-form fluxes (see Section 4). More concretely, the cancellation condition for the tadpole arising from the RR 4-form C_4 is (a D3-brane has f-flux quantum numbers $(n^i, m^i) = (1, 0)$)

$$2 \sum_a N_a n_a^1 n_a^2 n_a^3 = 32. \tag{7}$$

Furthermore, the cancellation conditions for the 8-form tadpoles yields:

$$\begin{aligned} 2 \sum_a N_a m_a^1 m_a^2 n_a^3 &= -32, \\ 2 \sum_a N_a m_a^1 m_a^3 n_a^2 &= -32, \\ 2 \sum_a N_a m_a^2 m_a^3 n_a^1 &= -32. \end{aligned} \tag{8}$$

The condition that all branes are mutually supersymmetric has the form:

$$\sum_i \arctan\left(\frac{\Im(T^i)}{f_a^i}\right) = 0. \tag{9}$$

This condition will fix some of the Kähler moduli T^i.

Finally, let us consider a configuration of 3 stacks of different D7-branes with spectrum identical to the MSSM, which is a concrete, type IIB mirror realization of the local brane set up in Fig. 1. Specifically, it is built by the f-flux quantum numbers [31] given in Table 2.

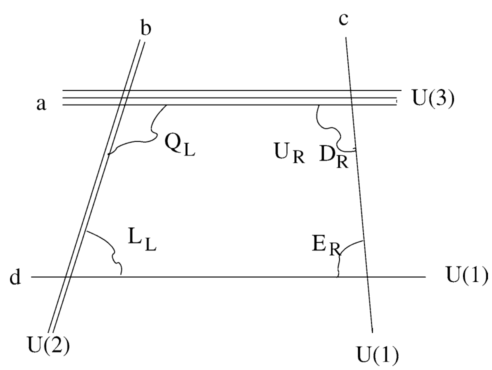

Fig. 1. Local D-brane set up for the Standard Model.

Table 2
MSSM D7-brane configuration with f-flux numbers (m^j, n^j)

Stack	Gauge group	(m^1, n^1)	(m^2, n^2)	(m^3, n^3)	N_a
1	$U(4)$	(0, 1)	$(1, g)$	$(-1, g)$	4
2	$SU(2)$	(1, 0)	(0, 1)	$(-1, 0)$	2
3	$SU(2)$	(1, 0)	$(-1, 0)$	(0, 1)	2

The corresponding gauge group is $G = U(4) \times U(2) \times U(2)$.[1] Only stack 1 carries non-trivial f-flux. For $g = 3$, there are 3 chiral generations of supersymmetric MSSM matter fields, namely lefthandend matter fields in the representations $3(\underline{4}, \underline{2}, \underline{1})$ from open strings stretching between the (12)-branes, 3 righthanded matter fields in the representations $3(\underline{4}, \underline{1}, \underline{2})$ from the (13) open string sector and a Higgs multiplet in the representations $(\underline{1}, \underline{2}, \underline{2})$ from the (23)-sector. By pulling apart the first stack of branes into a stack of $3 + 1$ D7-branes, the $SU(4)$ gauge group is Higgsed to $SU(3) \times U(1)_{B-L}$, and the matter fields decompose into the known SM representations of quarks and leptons. Then the four stacks of D7-brane precisely correspond to the brane set up in Fig. 2 in the T-dual type IIA picture.

Being supersymmetric, the 3 stacks of D-branes satisfy the supersymmetry conditions provided the 3 Kähler moduli obey the following two conditions:

$$T^2 = T^3, \qquad \arctan\big(\Im\big(T^2\big)/3\big) + \arctan\big(\Im\big(T^2\big)/4\big) = \pi/2 + \arctan\big(\Im\big(T^1\big)/2\big). \tag{10}$$

These 3 stacks of D7-branes will be a subsector in any concrete global model that satisfies the Ramond tadpole conditions by the addition of fluxes and some additional hidden sectors (for recent, concrete orbifold examples see [32–38]; for supersymmetric CFT constructions see [39–42]).

3. The effective low-energy action from open/closed string amplitudes

In $N = 1$ supersymmetric models the low-energy supergravity action of the massless gauge fields A_μ^I and matter fields Φ_a is determined by three moduli-dependent functions [43]:

- the gauge kinetic function $f_{IJ}(M)\, W_\alpha^I W_\alpha^J$;
- the superpotential $W = \hat{W}(M) + W_{abc}(M)\Phi_a\Phi_b\Phi_c + \cdots$;
- the Kähler potential $K = \hat{K}(M, \bar{M}) + K(M, \bar{M})_{a\bar{b}}\, \Phi_a\bar{\Phi}_b + \cdots$.

The matter field Kähler potential together with the cubic matter field superpotential is needed to compute the physical Yukawa couplings

$$Y_{abc} = e^{\hat{K}/2}\sqrt{K_{a\bar{d}}^{-1} K_{b\bar{e}}^{-1} K_{c\bar{f}}^{-1}}\, W_{def}. \tag{11}$$

We [10] will compute f, K, W of brane world models by calculating string tree level scattering amplitudes on the disk, which contain N_o open and N_c closed strings as in Fig. 2. The boundary of disk, being conformally equivalent to the upper half plane $\mathbb{H}_+$, is attached to the D-brane world volume. Hence we are using the following correlators on $\mathbb{H}_+$:

$$\begin{aligned} &\big\langle X^a(z_1)X^b(z_2)\big\rangle = -g^{ab}\log(z_1 - z_2), \qquad \big\langle X^a(z_1)X^b(\bar{z}_2)\big\rangle = -D^{ab}\log(z_1 - \bar{z}_2), \\ &\big\langle \psi^a(z_1)\psi^b(z_2)\big\rangle = \frac{g^{ab}}{z_1 - z_2}, \qquad \big\langle \psi^a(z_1)\bar{\psi}^b(\bar{z}_2)\big\rangle = \frac{D^{ab}}{z_1 - \bar{z}_2}. \end{aligned} \tag{12}$$

D^{ab} depend on the open string boundary conditions, i.e. on the F-flux:

$$D = -g^{-1} + 2(g + F)^{-1}. \tag{13}$$

Consider the following mixed open/closed string disk amplitude:

$$A_{N_c+N_o} \sim \big\langle V_c^1(\bar{z}_1, z_1)\cdots V_c^{N_c}(\bar{z}_{N_c}, z_{N_c})V_o^1(w_1)\cdots V_o^{N_o}(w_{N_o})\big\rangle. \tag{14}$$

[1] Note that in some orbifold models, the N_a will take values different from those in Table 2, if the D-branes are fixed under the orbifold group $Z_N \times Z_N$ and ΩI_n; e.g., for the $Z_2 \times Z_2$ orientifold, $N_1 = 8$ because the corresponding gauge group is broken to $U(N_1/2)$ by the orbifold symmetry.

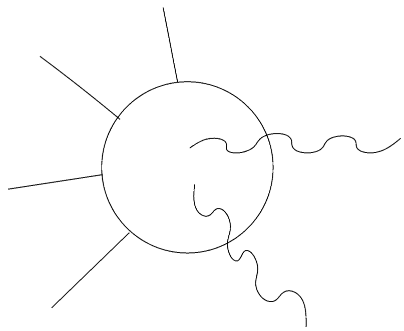

Fig. 2. The calculation model.

Here the open string vertex operators have the structure $V_o(z) = V(z)$ at the boundary $z = \bar{z}$; on the other hand, the closed string vertex operators live on the entire $\mathbb{H}_+$ and have the form $V_c(\bar{z}, z) = \bar{V}_c(\bar{z})V_c(z)$. The closed string vertex operators formally have two momenta, one along the D-brane, $k_\parallel$, and one transversal momentum, $k_\perp$, respectively $k = q$ and $k = Dq$ ($k_\parallel = (q + Dq)/2$). Therefore the above amplitude involves the integration over $2N_c + N_o - 3$ real positions z_i. This leads to a non-trivial momentum dependence already for the 3-point amplitude with $N_c = 1$, $N_o = 2$, namely A_3 is of $O(k^2)$. Furthermore correlators between holomorphic and anti-holomorphic operators in general contribute at the boundary. Hence there are less restrictions from internal charge conservation. In summary, these kind of mixed amplitudes are not just the square roots of the corresponding closed (heterotic) string matter amplitudes.

Now let us discuss a few specific amplitudes.

3.1. *The gauge kinetic function*

The gauge kinetic function for type IIB gauge fields on D3-branes and D7-branes with f-flux (or also for D5/D9-branes with f-flux) can be derived from the 3-point function between 2 gauge bosons and one modulus (Fig. 3). The explicit computation of the 3-point function yields [10]

$$A_3 = \langle A^{a_1} A^{a_2} T^j \rangle = \frac{i\bar{D}^j}{2T_2^j}\left[\frac{1}{2}(p_1 p_2)(\xi_1\xi_2) - \frac{1}{2}(p_1\xi_2)(p_2\xi_1)\right] t \frac{\Gamma(-2t)}{\Gamma(1-t)^2}, \quad t = p_1 p_3. \tag{15}$$

This amplitude is proportional to $\frac{\partial g_I^{-2}}{\partial M}$ ($M = S, T^i$). So by integrating above equation and taking into account the $N = 1$ SUSY condition,

$$\sum_{j=1}^{3} \frac{f^j}{\Im(T^j)} = \prod_{j=1}^{3} f^j \Im(T^j) \quad \bigl(f^j = b^j + 2\pi\alpha' F^j\bigr),$$

we derive for the gauge kinetic functions:

$$f_{\mathrm{D3}} = S, \qquad f_{\mathrm{D7}_i} = |m^k m^l| \bigl(T^i - f^k f^l S\bigr). \tag{16}$$

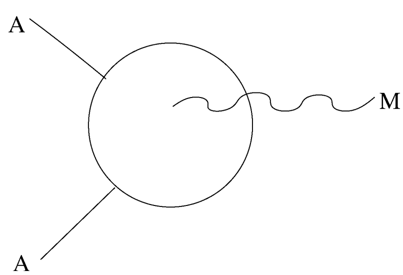

Fig. 3. 3-point function between 2 gauge bosons and one modulus.

Note that for D7-branes with f-flux, the gauge kinetic function depends both on S and on the Kähler moduli T^i. For a D7-brane without f-flux, the gauge kinetic function only depends on the modulus T^i of the torus, transversal to its world volume.[2]

3.2. The matter field Kähler potential

To compute the Kähler metrics of the matter fields (see also [44,45]) we need the 3-point amplitude shown in Fig. 4. It will give us information on the moduli dependence of the derivative of the Kähler metric: $A_3 \sim \partial K_{C^a\bar{C}^b}/\partial M^i$. In addition we need also the 4-point function between two moduli and two matter fields, as show in Fig. 5. This amplitude will allow us to get additional information about the Riemann tensor derived from the matter fields Kähler metric [46]:

$$A_4 \sim K_{C^a\bar{C}^b} K_{M^i\bar{M}^j} \frac{ts}{u} + s R_{C^a\bar{C}^b M^i\bar{M}^j}.$$

Let us first consider the Kähler metric for those matter fields which originate from open strings with endpoints on one single stack of D-branes. They transform in the adjoint representation of the gauge group G. In conformal field theory language they correspond to untwisted open string vertex operators. The metrics $G_{C_i^3\bar{C}_j^3}$ and $G_{C_i^{7,j}\bar{C}_i^{7j}}$ for the untwisted matter fields C_i^3 and C_i^{7j} may be obtained from the following differential equation, which follows from A_3:

$$\partial_{\Im T^j} G_{C_i\bar{C}_i} = \frac{D^j + \bar{D}^j}{4\Im T^j}\left(1 - 2\delta^{ij}\right) G_{C_i\bar{C}_i}. \tag{17}$$

Via integration we obtain for the D3-brane matter fields

$$G_{C_i^3\bar{C}_i^3} = \frac{\kappa_4^{-2}}{(U^i - \bar{U}^i)(T^i - \bar{T}^i)}, \quad i = 1, 2, 3. \tag{18}$$

The C_i^3 are the matter fields which describe the positions of the D3-branes in each of the three internal tori. They are the scalars of an $N = 4$ vector supermultiplet.

Let us now move on to the untwisted D7-matter fields $C_i^{7,j}$. For concreteness, let us consider the fields $C_i^{7,3}$, i.e. we shall discuss the case of a D7-brane, which is transversal to the third torus $T^{2,3}$. In this specific case, we find:

$$G_{C_1^{7,3}\bar{C}_1^{7,3}} = \frac{\kappa_4^{-2}}{(U^1 - \bar{U}^1)(T^2 - \bar{T}^2)} \frac{|1 + i\tilde{f}^2|}{|1 + i\tilde{f}^1|},$$

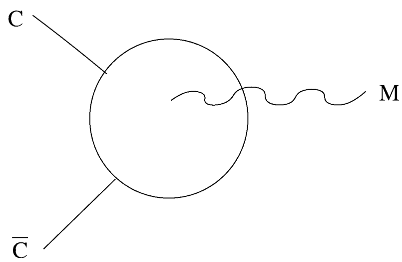

Fig. 4. 3-point amplitude for calculation of the Kähler metrics.

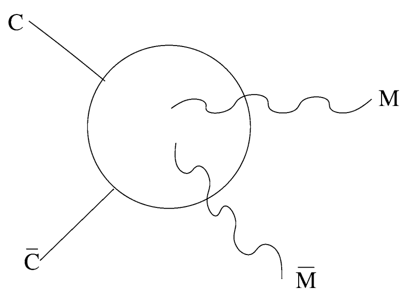

Fig. 5. 4-point function between two moduli and two matter fields.

[2] Note that there is a field redefinition involved going from the Kähler modulus in the string basis to the one in the supergravity basis.

$$G_{C_2^{7,3}\bar{C}_2^{7,3}} = \frac{\kappa_4^{-2}}{(U^2-\bar{U}^2)(T^1-\bar{T}^1)}\frac{|1+i\tilde{f}^1|}{|1+i\tilde{f}^2|},$$
$$G_{C_3^{7,3}\bar{C}_3^{7,3}} = \frac{\kappa_4^{-2}}{(U^3-\bar{U}^3)(S-\bar{S})}\left|1-\tilde{f}^1\tilde{f}^2\right|,$$
$$G_{C_i\bar{C}_k} = 0, \quad i \neq k, \tag{19}$$

where $\tilde{f}^i = \frac{f^i}{\Im T^i}$ is the physical 2-form flux. Here the indices $i = 1, 2$ refer to scalar fields living on the world volume of the D7-branes, constituting an $N = 2$ hypermultiplet, whereas the scalar $C_3^{7,3}$ describes the position of the D7-brane on $T^{2,3}$, being member of an $N = 2$ vector multiplet. The other cases $G_{C_i^{7,j}\bar{C}_i^{7,j}}$ with $j = 1, 2$ are obtained from the above results by permuting fields.

Now let us come to those matter fields which located at the intersection of two D6-branes at angle θ respectively between D3–D7$_a$, D7$_a$–D7$_b$ respectively D5–D9$_a$, D9$_a$–D9$_b$ branes with fluxes F_θ. These are the fields, which transform in bifundamental representations, and which play the role of the Standard Model matter fields in any realistic model. From the open string conformal field theory point of view these fields correspond to twisted vertex operators, since the open strings end on different stacks of D-branes. In case the two intersecting branes form a supersymmetric 1/4 BPS configuration, i.e. the D7-branes wrap different 4-cycles, the 3-point function takes the form:

$$A_3 = \langle C^\theta T^j \bar{C}^{-\theta}\rangle \sim -e^{-i\pi\theta^j}\frac{\Gamma(-2t)[\Gamma(-t+\theta^j)+(1-\theta^j)\Gamma(-1-t+\theta^j)]}{\Gamma(-1-t+\theta^j)\Gamma(-t+\theta^j)\Gamma(1-t-\theta^j)}. \tag{20}$$

From that we obtain the following differential equation:

$$\frac{\partial K_{C_\theta\bar{C}_\theta}}{\partial T_j} \sim e^{-\pi i\theta_j}\sin(\pi\theta_j)\left[2\gamma_E+\psi(\theta_j)+\psi(1-\theta_j)\right]K_{C_\theta\bar{C}_\theta}. \tag{21}$$

Integrating this equation and taking the limit $\alpha' \to 0$ yields:

$$G_{C^{7_a7_b}\bar{C}^{7_a7_b}} \sim \prod_{j=1}^{3}\sqrt{\frac{\Gamma(\theta_{ab}^j)}{\Gamma(1-\theta_{ab}^j)}}. \tag{22}$$

Here the angle θ_{ab}^j, reminiscent from the intersecting D6-brane description, encodes the two flux components on the different stacks a and b of D-branes:

$$\theta_{ab}^j = \frac{1}{\pi}\left[\arctan\left(\frac{f_b^j}{\Im(T^j)}\right) - \arctan\left(\frac{f_a^j}{\Im(T^j)}\right)\right]. \tag{23}$$

Further moduli dependences can be derived from the 4-point amplitude $A_4 = \langle C^\theta U^m \bar{U}^m \bar{C}^{-\theta}\rangle$. Then we obtain the following final result for the twisted (1/4 BPS) matter field Kähler metric [10]:

$$G_{C^{7_a7_b}\bar{C}^{7_a7_b}} = \kappa_4^{-2}\prod_{j=1}^{3}\left(U^j-\bar{U}^j\right)^{-\theta_{ab}^j}\sqrt{\frac{\Gamma(\theta_{ab}^j)}{\Gamma(1-\theta_{ab}^j)}}. \tag{24}$$

By similar methods we can also derive the Kähler metric for two intersection branes that are a 1/2 BPS configuration, e.g., two D7-branes which are transversal with respect to the same torus:

$$G_{C^{7_27_3}\bar{C}^{7_27_3}} = \frac{\kappa_4^{-2}}{(S-\bar{S})^{1/2}(T^1-\bar{T}^1)^{1/2}}\frac{1}{(U^2-\bar{U}^2)^{1/2}(U^3-\bar{U}^3)^{1/2}}. \tag{25}$$

Finally let us also give the result for the Kähler potential of the bulk moduli fields, which can be derived from scattering of the corresponding closed string states on the D-branes:

$$\kappa_4^2\hat{K} = -\ln(S-\bar{S}) - \sum_{i=1}^{3}\ln(T^i-\bar{T}^i) - \sum_{i=1}^{3}\ln(U^i-\bar{U}^i). \tag{26}$$

3.3. *The matter field superpotential/Yukawa couplings*

At the end of this section we discuss very briefly the computation of the twisted matter field Yukawa couplings [47,48,10] and the associated cubic matter field superpotential. Now we need to compute the 4-point amplitude $\langle C_\theta \bar{C}_\theta C_\theta \bar{C}_\theta \rangle$. Taking a suitable factorization limit this amplitude can be used to derive the following expression for the physical Yukawa coupling among 3 twisted chiral matter fields

$$Y_{j,k,-j-k} \sim \prod_{i=1}^{3} \left[\frac{\Gamma(1-\theta_j)\Gamma(1-\theta_k)\Gamma(\theta_j+\theta_k)}{\Gamma(\theta_j)\Gamma(\theta_k)\Gamma(1-\theta_j-\theta_k)} \right]^{1/4} W_{j,k,-j-k}. \tag{27}$$

W is the exponential superpotential describing the classical world sheet instantons: $W_{j,k,-j-k} \sim e^{-A_{j,k,-j-k}}$, A being the area between the 3 D6-branes in type IIA. Note that using the general expression Eq. (11) for the Yukawa couplings, above equation is consistent with the previous result for the twisted matter field kinetic energies in Eq. (24).

4. The 3-form flux induced soft SUSY breaking terms

Now we will study the effect of turning on non-vanishing bulk 3-form flux in type IIB orientifolds [13–21]:

$$\frac{1}{(2\pi)^2\alpha'} \int_{C_3} G_3 \neq 0, \qquad G_3 = F_3 - SH_3, \qquad N_{\text{flux}} = \frac{1}{(2\pi)^4\alpha'^2} \int H_3 \wedge F_3. \tag{28}$$

Non-vanishing 3-form fluxes have several interesting effects:

– A modified Ramond tadpole condition:

$$N_{\text{flux}} + 2\sum_a N_a n_a^1 n_a^2 n_a^3 = 32. \tag{29}$$

– A modified D-term scalar potential:

$$V_D(T^i) = V_{\text{D3/D7}} + V_{\text{O3/O7}} - e^{-\phi_{10}} T_3 N_{\text{flux}}. \tag{30}$$

As before, V_D fixes some of the Kähler moduli T^i.

– The generation of a superpotential, which depends on the dilaton S and on the complex structure moduli U^i, but not on the Kähler moduli T^i:

$$\hat{W}(S, U^i) = \frac{1}{(2\pi)^2\alpha'} \int G_3 \wedge \Omega. \tag{31}$$

The associated scalar potential has the form:

$$\hat{V}_F(S, U^i) = \frac{1}{(2\pi)^7\alpha'^4} \int G^{\text{IASD}} \wedge \star_6 \bar{G}^{\text{IASD}}, \quad \star_6 G^{\text{(A)ISD}} = \pm G^{\text{(A)ISD}}. \tag{32}$$

Possible flux vacua are defined by minimization of the scalar potential:

$$\frac{\partial \hat{V}_F}{\partial S} = \frac{\partial \hat{V}_F}{\partial U^i} = 0. \tag{33}$$

So minimizing $\hat{V}_F$, which is of the standard supergravity form, $\hat{V}_F = \hat{K}_{ij} F^i \bar{F}^j - 3e^{\hat{K}} |\hat{W}|^2$, generically fixes the complex structure moduli U^i and the dilaton S, but leaves the Kähler moduli undetermined (these are partially fixed by the D-term potential V_D). Concerning the breakdown of $N = 1$ supersymmetry there are possible classes of vacua:

(i) *Supersymmetric vacua*: here all F-terms vanish in a given vacuum solution: $F^{T^i} = F^{U^i} = F^S = 0$. One can show that in this the 3-form flux must be a primitive ISD (2,1)-form: $G_{2,1} \neq 0$.

(ii) *Non-supersymmetric vacua*: here $N = 1$ supersymmetry is spontaneously broken by one or all non-vanishing F-terms. More specifically, there are the following 3 possibilities:
  - $F^T = D_{\bar{T}} \hat{W} = \kappa_4^2 \hat{W} \partial_T \hat{K} \sim \int G_3 \wedge \Omega \neq 0$. In this case G is a ISD (0,3)-form: $G_{0,3} \neq 0$.
  - $F^S = D_{\bar{S}} \hat{W} = \partial_S \hat{W} + \kappa_4^2 \hat{W} \partial_S \hat{K} \sim \int \bar{G}_3 \wedge \Omega \neq 0$. In this case G is a IASD (3,0)-form: $G_{3,0} \neq 0$.
  - $F^{U^i} = D_{\bar{U}^i} \hat{W} = \partial_{U^i} \hat{W} + \kappa_4^2 \hat{W} \partial_{U^i} \hat{K} \neq 0$. In this case G is a IASD (1,2)-form: $G_{1,2} \neq 0$.

To be more specific let us consider 3-form fluxes in toroidal and orbifold compactifications. In fact, not all flux components in will survive the orbifold projection. In addition, some or all complex moduli U^i will be frozen to discrete values by the $Z_N \times Z_M$ modding. In the $Z_2 \times Z_2$ orbifold, 8 complex flux components indeed survive and also all of the U^i remain unfixed. However, e.g., for the Z_3 orbifold, only $G_{(3,0)}$ and $G_{(0,3)}$ are allowed, and all U^i are frozen to $U^i = 1/2 + i\sqrt{3}/2$. As a general result of this investigation it turns out that only the IASD flux $G_{(3,0)}$ and the ISD flux $G_{(0,3)}$ are generic for all orbifolds. Therefore we will concentrate in the following discussion on these two complex fluxes. Expressed in terms of a complex basis, the $G_{(3,0)}$ and $G_{(0,3)}$ fluxes take the following form:

$$\frac{1}{(2\pi)^2\alpha'}G_{03} = A_0\big(d\bar{z}^1 \wedge d\bar{z}^2 \wedge d\bar{z}^3\big),$$
$$\frac{1}{(2\pi)^2\alpha'}G_{30} = B_0\big(dz^1 \wedge dz^2 \wedge dz^3\big). \tag{34}$$

Let us just remark that due to the absence of the ISD (2, 1) 3-form fluxes in many of the orbifold models, supersymmetric flux solutions do not exist.

In orbifold compactifications, we obtain for the superpotential:

$$\hat{W} = \big(a^0 - Sc^0\big)U^1U^2U^3 - \big\{\big(a^1 - Sc^1\big)U^2U^3 + \big(a^2 - Sc^2\big)U^1U^3 + \big(a^3 - Sc^3\big)U^1U^2\big\}$$
$$- \sum_{i=1}^{3}(b_i - Sd_i)U^i - (b_0 - Sd_0), \tag{35}$$

where the a^i, b^i, c^i, d^i are real, integer flux coefficients. Then the explicit expressions for F^T and F^S are the following:

$$\bar{F}^{\bar{S}} = \big(S - \bar{S}\big)^{1/2}\big(T - \bar{T}\big)^{-3/2}\prod_{i=1}^{3}\big(U^i - \bar{U}^i\big)^{-1/2}\kappa_4^2\frac{\lambda}{(2\pi)^2\alpha'}\int \bar{G}_3 \wedge \Omega$$
$$= \lambda\kappa_4^2\big(S - \bar{S}\big)^{1/2}\big(T - \bar{T}\big)^{-3/2}\prod_{i=1}^{3}\big(U^i - \bar{U}^i\big)^{-1/2}$$
$$\times \Bigg\{\big(a^0 - \bar{S}c^0\big)U^1U^2U^3 - \big[\big(a^1 - \bar{S}c^1\big)U^2U^3 + \big(a^2 - \bar{S}c^2\big)U^1U^3 + \big(a^3 - \bar{S}c^3\big)U^1U^2\big]$$
$$- \sum_{i=1}^{3}\big(b_i - \bar{S}d_i\big)U^i - \big(b_0 - \bar{S}d_0\big)\Bigg\},$$
$$\bar{F}^{\bar{T}} = \big(S - \bar{S}\big)^{-1/2}\big(T - \bar{T}\big)^{-1/2}\prod_{i=1}^{3}\big(U^i - \bar{U}^i\big)^{-1/2}\kappa_4^2\hat{W}. \tag{36}$$

It is easy to see that $\hat{V}_F$ is only non-vanishing if the 3-form flux is IASD, i.e. for $G_{(3,0)}$ (and also for $G_{(1,2)}$). Specifically, when we express $G_{(3,0)}$ by its complex coefficient B_0, $\hat{V}_F$ becomes:

$$\hat{V}_F = \kappa_4^2 B_0^2 \frac{\prod_{i=1}^{3}|U^i - \bar{U}^i|}{|S - \bar{S}|\prod_{i=1}^{3}|T^i - \bar{T}^i|}. \tag{37}$$

Now we can combine the 3-form flux induced bulk effective action with the effective action for the matter fields on D3/D7-branes with f-flux. This will lead to soft supersymmetry breaking terms for the matter fields in case supersymmetry is spontaneously broken [22–26]. The general supergravity expressions for the soft scalar masses and the trilinear couplings are:

$$m^2_{a\bar{b},\mathrm{soft}} = \big(|m_{3/2}|^2 + \hat{V}_F\big)K_{C_a\bar{C}_b} - F^i F^{\bar{j}} R_{i\bar{j}a\bar{b}},$$
$$A_{abc} = F^i D_j\big(e^{\hat{K}/2}Y_{abc}\big). \tag{38}$$

For the soft scalar masses some simplifications are possible by plugging the actual expressions for F^S and F^T into above equation. For this and for applying the formulas to the MSSM construction at the end of Section 2, we refer to [24,49].

The gaugino masses have the following general form:

$$m_{gI} = F^i \partial_i \log(\mathrm{Im}\, f_I). \tag{39}$$

More specifically, for the gaugino masses for the gauge fields living on the D7-branes wrapped around the 4-cycle $T^{2,k} \times T^{2,l}$, using Eq. (16), we obtain:

$$m_{g,\mathrm{D7}_i} = F^S \frac{-f^k f^l}{(T^j - \bar{T}^j) - f^k f^l (S - \bar{S})} + F^{T^j} \frac{1}{(T^j - \bar{T}^j) - f^k f^l (S - \bar{S})}. \tag{40}$$

Finally, for the MSSM D7-brane set-up with wrapping number as specified in Table 2, this formula looks even simpler [49]:

$$m_{g,1} = \frac{F^{T^1} + g^2 F^S}{(T^1 - \bar{T}^1) + g^2 (S - \bar{S})} \quad \text{for stack } 1,$$
$$m_{g,j} = \frac{F^{T^j}}{(T^j - \bar{T} j1)} \quad \text{for stacks } 2, 3. \tag{41}$$

5. Gaugino condensation and 3-form fluxes

The last section steps a little bit aside from the discussion on intersecting branes and models D3/D7-branes with chiral fermions. We will rather discuss the generation of 3-form flux and its relation to gaugino condensation in the context of the AdS/CFT correspondence. Specifically we would like to pose the question whether certain 3-form fluxes correspond to the dynamical formation of a gaugino condensate in the gauge sector of the theory? The motivation to assume such a correspondence is the following:

– In the context of the large N transition (geometric transition) the flux is related to a gaugino condensate $\langle\lambda\lambda\rangle$ [50].
– In heterotic string compactifications with gaugino condensate one needs $(3,0)$ H-flux for stabilization of the vacuum.
– As we discussed before, in type IIB orientifolds $G_{3,0}$ is proportional to $F^S \neq 0$; this corresponds to the dilaton dominated spontaneous SUSY breaking, which is indeed very similar to gaugino condensation.

However from now on we are switching from compact models to non-compact supergravity backgrounds with 3-form fluxes. Via the AdS/CFT correspondence these are dual to certain supersymmetric gauge theory. As we will show [27] turning on a $(3,0)$ G-flux in the supergravity background indeed corresponds to the formation of a gaugino condensate in the dual globally $N=1$ supersymmetric field theory. Since the gaugino condensate does not spontaneously break global $N=1$ supersymmetry, the $(3,0)$-flux must be also a supersymmetric solution in type IIB supergravity on a non-compact background, in contrast to the compact case. This will be further discussed at the end of this chapter.

Specifically we will use the AdS/CFT correspondence between the $N=1^*$ $SU(N)$ gauge theory and Polchinski–Strassler solution [28] of type IIB supergravity with 3-form flux. The $N=1^*$ gauge theory is defined by a mass deformation of $N=4$ $SU(N)$ super-Yang–Mills theory, namely by turning on the following superpotential:

$$W = m_{ij}\phi_{(i}\phi_{j)}, \quad m_{ij} = m, \quad i = 1, 2, 3. \tag{42}$$

W explicitly breaks $N=4$ supersymmetry down to $N=1$ supersymmetry giving a common mass to all three chiral, adjoint supermultiplets inside the $N=4$ vector multiplet. So in the infrared, for mass scales much smaller than m, one is dealing with pure $N=1$ $SU(n)$ gauge theory; this theory possesses in the confining phase a non-vanishing gaugino condensate $\langle\bar{\lambda}\bar{\lambda}\rangle = m^3$ at large 'tHooft coupling $g^2_{\mathrm{YM}}N$. The dual supergravity solution, describing the RG flow from $N=4$ to $N=1$, has the form of a warped background with 3-form flux

$$ds^2 = Z^{-1/2}\eta_{\mu\nu}\,dx^{\mu\nu} + Z^{1/2}g_{mn}\,dx^m\,dx^n,$$
$$G_3 = F_3 - \tau H_3, \quad F_3 = dC_2, \quad H_3 = dB_2, \quad G_3 = G_3\big(x^m\big), \quad \tau = \tau\big(x^m\big). \tag{43}$$

However the exact supergravity solution is not known, but there exists only an iterative expansion in powers of m. At linear order in m with constant τ it was shown by [28] that turning on the 3-form flux in supergravity corresponds to turning on the mass parameter m in the gauge theory. This was extended by [51] to a solution quadratic in m and non-constant in τ. Without going into any detail, one can show [27] that at cubic order in m one gets

$$G^{(3,0)} = \partial A_2^{(2,0)}, \quad A_2^{(2,0)} = m^3 r^{-\Delta}\varepsilon_{ijk}\left(\frac{\bar{z}^2}{r^2}\right)^2 \frac{z^i}{r} d\left(\frac{z^j}{r}\right) \wedge d\left(\frac{z^j}{r}\right). \tag{44}$$

As it is known in AdS/CFT, this $A_2^{(2,0)}$ corresponds to $\langle\lambda\lambda\rangle$!

As already emphasized, $G^{(3,0)}$ preserves $N = 1$ supersymmetry. How can that be? The answer lies in the form how supersymmetry is realized on the supergravity background. Specifically, there exist several kinds of known supersymmetric Killing spinor solutions:

– Type A(ndy): This is realized in heterotic/type II strings with H-flux [52].
– Type B(ecker): its solution is the supersymmetric ISD $G_{2,1}$-flux in type IIB, already discussed in the section on G-flux compactifications [53–55]. The corresponding spinor ansatz for the above two cases reads

$$\epsilon(x, y) = a(y)\varepsilon(x) \otimes \eta_-(y) + b(y)\varepsilon^*(x) \otimes \eta_+(y), \tag{45}$$

where a and b denote complex functions, ε is the four-dimensional supersymmetry parameter and $\eta_+ = (\eta_-)^*$ is a globally defined spinor normalised to one. The existence of one globally defined spinor η implies that the tangent bundle over the transverse 6-dimensional space has an $SU(3)$ group structure. This ansatz includes the type A ansatz, where $b = a^*$, the type B, where $b = 0$, and the more general type discussed in [56], called type C.
– Type D(all'Agata): There exists an even more general supersymmetry ansatz [57], which can be expressed:

$$\epsilon(x, y) = a(y)\varepsilon(x) \otimes \eta_-(y) + \varepsilon^*(x) \otimes \big(b(y)\eta_+(y) + c(y)\chi_+(y)\big). \tag{46}$$

It is based on the existence of two globally defined spinors, η and χ, which are linearly independent. It implies that the group structure of the tangent bundle of the transverse 6-dimensional space is further reduced to $SU(2)$. In particular, using the ansatz (46), the Hodge type of the 3-form flux is no more constrained by supersymmetry, and one can now have $(3, 0)$ as well as $(0, 3)$ fluxes. Moreover, and in contrast with the case of $SU(3)$-structures, there is no preferred choice for the (almost) complex structure J, but one actually has a $U(1)$-worth of possibilities. Indeed, after some algebra one can show that the Polchinski/Strassler background at order m^3 with IASD flux $G^{(3,0)}$ satisfies the type D Killing spinor ansatz and is supersymmetric with an $SU(2)$ group structure. Let us finally remark that the type D ansatz is a manifestation of the dielectric nature of solution, namely that the underlying D3-branes are dissolved into D5-branes by the 3-form flux (see also [58]).

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