D-BRANE MODEL BUILDING, PART III

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PITP School, IAS Princeton, July 2008

VI- NEW COUPLINGS FROM STRINGY INSTANTONS

U(1) global symmetries are not forever

- We saw how in generic IIA and IIB D-brane models some $U(1)$'s (anomalous or not) can get a mass combining with some RR fields.
- The corresponding $U(1)_X$ symmetry remains perturbatively to all orders as an effective global symmetry. E.g. any superpotential operator

$$\frac{1}{M_s^{n-3}}\Phi_{q_1}...\Phi_{q_n} = 0 \; ; \; \sum_i q_i \neq 0 \tag{1}$$

- That is the case e.g. baryon number and lepton number in MSSM-like models of the type we have described. That could be a problem for baryo- or lepto-genesis and/or the existence of Majorana neutrino masses.
- It has been recently realized however that non-perturbative stringy instanton effects may give rise to superpotential operators

$$\frac{1}{M_s^{n-3}} e^{-M} \Phi_{q_1} \ldots \Phi_{q_n} \neq 0 \;;\; \sum_i q_i \neq 0 \quad (2)$$

with M a IIA complex structure (IIB Kahler modulus) field whose imaginary part shifts under a $U(1)_X$ gauge transformation of parameter Λ_x like

$$M \longrightarrow M + \Lambda_x \left(\sum_i^n q_i\right) \quad (3)$$

- This shift is such that the the operator is fully gauge invariant.

Majorana neutrino masses

- Let us recall some well known facts. The simplest explanation for the smallness of neutrino masses is the celebrated see-saw mechanism.

- If there are right-handed neutrinos ν_R^a with large Majorana masses M_M and standard Dirac masses M_D, the lightest eigenvalues have masses of order

$$M_\nu \simeq \frac{M_D^2}{M_M} , \tag{4}$$

of order experimental results for M_D of order of standard charged lepton masses and for $M_M \propto 10^{10} - 10^{13}$ GeV.

- Why ν_R exists? They are natural in left-right symmetric extensions of the SM like $SO(10)$, $SU(3) \times SU(2)_L \times SU(2)_R \times U(1)_{B-L}$.

- The ν_R may get Majorana masses through Yukawa couplings

$$\nu_R \nu_R \phi_M \tag{5}$$

In $SO(10)$ one has $\phi_M = 126$, in $3 \times 2 \times 2 \times 1$ it is $(1, 1, 3, 2)$.

- Or else through some non-renormalizable couplings

$$\frac{1}{M_{B-L}} \nu_R \nu_R < \bar{N}_R \bar{N}_R > \tag{6}$$

- In the first case in the MSSM R-parity is automatic but the Higgs sector is ugly..... In the second case the Higgs sector is simple but R-parity is not automatic.....

Some general features for ν_R masses in string models

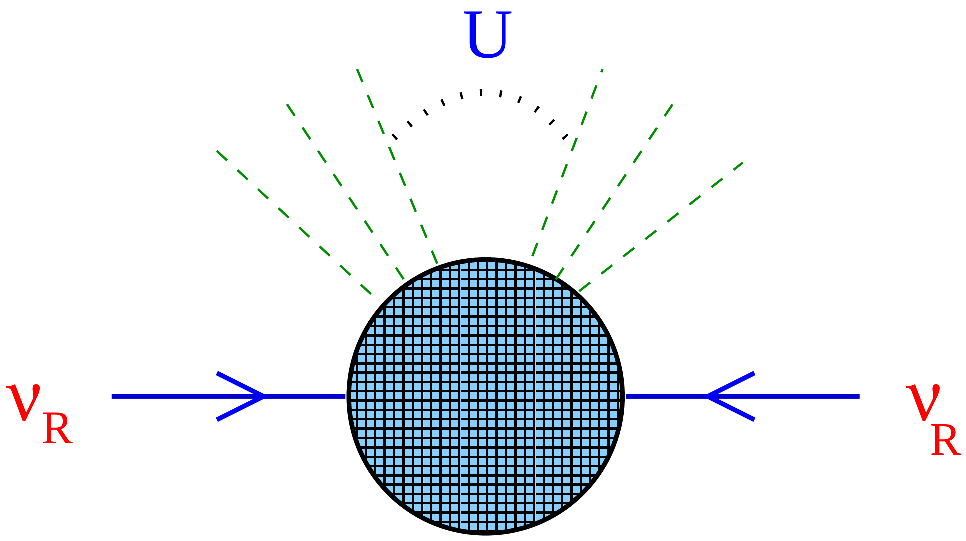
- Note that unlike field theory models, string vacua are quite rigid. Given a compactification the spectrum and couplings are fixed (as a function of the moduli). We cannot 'complete our model' with the fields and couplings we would like to have!!
- The presence of ν_R's is generic in SM-like vacua.
- Dirac neutrino masses are generically present but couplings generating Majorana ν_R-masses are absent!!. The required Higgs fields have 'charge too big' to be in the massless spectrum.
- Practically all MSSM-like models constructed to date have a $U(1)_{B-L}$ gauge boson beyond the SM group.
- Efforts up to now in order to obtain ν_R masses in string models typically resort to higher-dimensional $\nu_R \nu_R \phi^n$ operators. But then typically one gets 1) Too small ν_R masses 2) R-parity violation and 3) Often fast proton decay.

- In D-brane models B and L symmetries are $U(1)$ gauged symmetries. Then right-handed neutrinos ν_R are perturbatively massless.
- Even if $U(1)_{B-L}$ gets massive by combining with a RR-field, it will remain perturbatively massless.
- However String Theory instantons may generate ν_R-masses through operators [a]

$$e^{-U}\ \nu_R \nu_R M_{string} \tag{7}$$

with U axion-like fields which shift under $U(1)_{B-L}$.

[a] L.E.I. and A. Uranga, hep-th/0609213; R. Blumenhagen, M. Cvetic, T. Weigand hep-th/0609191.



- This happens only in a restricted class of compactifications in which $U(1)_{B-L}$ is massive through a Stuckelberg term (like in GS mechanism).
- Such instantons respect R-parity in MSSM-like models.
- String instantons may generate other interesting superpotentials like e.g. a μ-term or some perturbatively forbidden Yukawas .

  We will ilustrate the mechanism in terms of the intersecting D6-brane approach but it should be clear the mechanism is GENERAL in String Theory.

String Theory Instantons

- We saw in Type IIA orientifolds gauge groups live on D6-branes wrapping 3-cycles in the CY.
- We now consider the effect of D2 euclidean branes also wrapping a 3-cycle. They have Drichlet b.c. in Minkowski, localized in space and time, instantons [a].
- They may contribute in semiclassical tunneling processes, as standard gauge theory instantons do.
- Standard gauge theory instantons would correspond to D2-branes wrapping precissely the same 3-cycle as the D6-branes where the gauge group lives.
- However, in string theory there are other instanton varieties corresponding to D2-branes wrapping other cycles and intersecting the D6-branes present.
- At the D2-D6 brane intersections live fermionic zero modes which are

[a] Becker^2,Strominger (95);Witten (96,99);Harvey,Moore (99)

charged under the D6 gauge group.

- They may give rise to superpotentials involving 4-D matter fields [b]. The kind of operator which is interesting for as is:

$$e^{-S_{D2}}\,\nu_R\nu_R \tag{8}$$

- The operator $\nu_R\nu_R$ has charge = 2 under both $U(1)_{B-L}$ and $U(1)_R$. Thus the operator $e^{-S_{D2}}$ has to transform with charges = -2.

- Transition amplitude induced by D2-instanton M is proportional to (S_{D^2}= Born-Infeld action)

$$e^{-S_{D2}} \;=\; \exp(\,-\frac{V_{\Pi_M}}{\lambda}\,+\,i\sum_r q_{M,r}a_r\,) \tag{9}$$

(e.g. in toroidal models $q_M = n_M m_M m_M, n_M n_M n_M$ in terms of wrapping numbers of $D2$)

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[b] Ganor (96);Florea,Kachru,Mc Greevy,Saulina (06)

- We already saw that the $U(1)_A$ gauge bosons have non-trivial couplings to a set of basic 2-forms B_r in the 4d theory

$$S_{BF} = \sum_{A,r} N_A p_{Ar} \int_{4d} B_r \wedge F_A \tag{10}$$

(e.g. in toroidal models $p_A = m_A n_A n_A$ or $m_A m_A m_A$)

- This implies that under a $U(1)_A$ gauge transformation :
$A_A \to A_A + d\Lambda_A$ the a_r scalar dual to B_r transforms:

$$a_r \to a_r + N_A (p_{Ar} - p_{A^*r})\Lambda_A \tag{11}$$

$$\sum_r q_{M,r} \sum_A N_A (p_{Ar} - p_{A^*r})\Lambda_A = \sum_A N_A (I_{MA} - I_{MA^*})\Lambda_A \tag{12}$$

where $I_{MA} = \Pi_M \cdot \Pi_A$ is the intersection number of M and A,

- Then

$$e^{-S_{D2}} \longrightarrow \exp\left(-i\sum_A N_A\,(\,I_{MA} - I_{MA^*}\,)\,\Lambda_A\,\right) e^{-S_{D2}} \qquad (13)$$

- In the case of real instantons $M = M^*$ one rather has

$$e^{-S_{D2}} \longrightarrow \exp\left(-i\sum_A N_A\,(\,I_{MA}\,)\,\Lambda_A\,\right) e^{-S_{D2}} \qquad (14)$$

- In both cases in order that $e^{-S_{D2}}$ has $U(1)_{B-L}, U(1)_R$ charge = -2 the instanton M must obey[a]:

$$I_{Ma} = I_{Ma^*} = I_{Mb} = I_{Mb^*} = 0 \; ; \; I_{Mc} - I_{Mc^*} = I_{Md} - I_{Md^*} = 2 \qquad (15)$$

[a] L.E.I. and Uranga (06); Blumenhagen,Cvetic and Weigand (06)

D2 instantons and fermionic zero modes

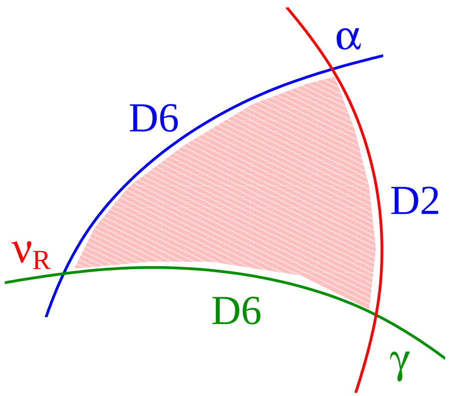
- The computation of the effective action involves an integration over the fermionic zero modes of the instanton.
- There are two types of fermionic zero modes:
  - Zero modes NOT charged under D=4, D6 gauge group
  - Zero modes charged under D=4, D6 gauge group
- There must be exactly two uncharged fermion zero modes in order to saturate the $d\theta^2$ integration in the effective action [a].
- That is in principle only possible for D2 branes with Chan-Paton symmetries $O(1)$ ($M = M^*$).
- However it has been shown that instantons with $U(1)$ or $Sp(2)$symmetries may also be effective in the presence of certain multi-instanton effects which can give masses to unwanted extra zero modes.

[a] Witten (96)

The Microscopic Mechanism

- How such a term is generated? There are open strings stretching between D2 and the background D6.
- Quantization of these open strings shows there are fermionic zero modes α_i, γ_i i=1,2 at the Mc and Md* intersections respectively.
- There are Mc-cd*-d*M cubic couplings involving the scalar ν_R:

$$L_{cubic} \propto d_a^{ij} \left(\alpha_i \, \nu^a \gamma_j\right) \; , a = 1, 2, 3 \tag{16}$$



- We have to integrate over zero modes α_i, γ_i : (recall e.g. $\int d\alpha\alpha = 1, \int d\alpha = 0$)

$$\int d^2\alpha\, d^2\gamma\, e^{-d_a^{ij}\,(\alpha_i \nu^a \gamma_j)} \propto -\nu_a\nu_b \int d^2\alpha\, d^2\gamma\, \alpha_i\alpha_j\gamma_k\gamma_l\, d_a^{ik} d_b^{jl}$$
$$= \nu_a\nu_b\,(\,\epsilon_{ij}\epsilon_{kl} d_a^{ik} d_b^{jl}\,) \quad (17)$$

yielding a right-handed neutrino mass term:

$$\nu_a\nu_b M_s(\,\epsilon_{ij}\epsilon_{kl} d_a^{ik} d_b^{jl}\,)\exp(-\frac{V_{\Pi_M}}{\lambda} + i\sum_r q_{M,r}a_r\,) \quad (18)$$

- The gauge $U(1)_c$,$U(1)_d$ transformation of the bilinear piece and the $e^{-S_{D2}}$ factor nicely cancel.
- Note that the flavor structure is controled by the d_a^{ik} coefficients which are given by the triangle correlators.

Size of ν_R Majorana masses

- The ν_R Majorana masses are of order:

$$M_{\nu_R} \simeq M_s\, d^2 \exp(-\frac{V_{\Pi_M}}{\lambda}) = M_s\, d^2 \exp(-\sum_r q_{M,r} ReU_r) \qquad (19)$$

- In the $N=1$ SUSY case U_r are the complex structure moduli.
- For usual gauge instantons this would be strongly supressed since $\sum_r q_{M,r} ReU_r$ is nothing but the inverse gauge coupling constant.
- For the non-gauge instantons here considered that quantity is unrelated to SM gauge couplings, may be of order one.
- One expects ν_R masses to be e.g. a few orders of magnitude below the string scale M_s.

R-parity and instanton induced Majorana masses

- The composite field $\exp(-U)$ has $U(1)_{B-L}$ and $U(1)_R$ charges = -2. It is a sort of effective Majoron-like field.
- This means that there is an unbroken gauge Z_2 soubgroup of $U(1)_{B-L}$.
- This is nothing but R-parity which is an automatic symmetry if ν_R masses are generated this way.
- Of course this is so in the absence of other possible instantons which might violate R-parity. One has to check in each compactification for the absence of these other instantons.

Further remarks

- Note that it is possible that in the preseence of RR and/or NS fluxes many or all extra zero modes may be lifted. (E.g. that happens in IIB orientifolds with 3-form fluxes in wich certain D3-instanton zero modes are lifted). In that case superpotentials may be generated event if apparently there are too many zero modes.

- Other class of extra zero modes may arise if the instanton M intersects 'hidden sector' D6-branes. In such a case operators of the general form

$$\nu_a \nu_b \Phi_H ....\Phi_H \tag{20}$$

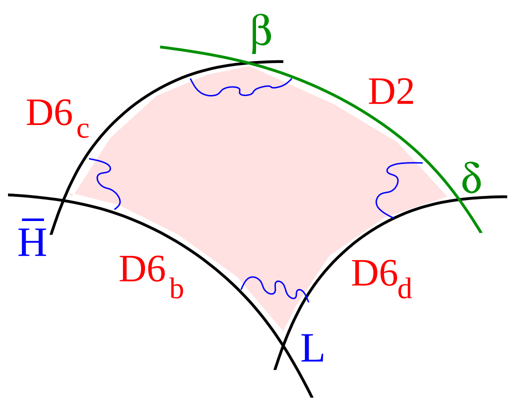
may appear, with Φ_H hidden sector D=4 fields . These may give rise tu ν_R masses upon Φ_H vevs.

- Note that although our discussion uses the language of $N = 1$ SUSY the mechanism is general and should exist also in non-SUSY models.

Direct generation of Weinberg operator

- In addition to the see-saw contribution there might be a direct generation of the Weinberg operator $L\bar{H}L\bar{H}$.
- Instantons W with intersections

$$I_{Wc} - I_{Wc^*} = I_{Wd} - I_{Wd^*} \; = \; -2 \tag{21}$$



- A Weinberg operator is generated:

$$L_a \bar{H} L_b \bar{H} \frac{1}{M_s} (\, \epsilon_{ij}\epsilon_{kl} \sum_s c_a^{ik}(s) c_b^{jl}(s)\,) \exp(\,-U_s) \tag{22}$$

- This contribution may be comparable to the see-saw one.

Neutrino Mass Matrices : the Weinberg Operator

- The flavour structure is particularly simple for $Sp(2)$ instantons [a] In that case the fermionic zero modes are $Sp(2)$ doublets and

$$c_a^{ij} = \epsilon^{ij} c_a \; ; \; d_a^{ij} = \epsilon^{ij} d_a \quad , \quad a = 1, 2, 3 \tag{23}$$

- Then the left-handed neutrino masses from Weinberg operators are:

$$M^{\nu_L}{}_{ab} = \frac{2 < \overline{H} >^2}{M_s} \sum_r c_a^{(r)} c_b^{(r)} \, e^{-U_r} \tag{24}$$

- The sum is over different instanton contributions. One thus has a structure

$$M^{\nu_L} = \frac{2V^2}{M_s} \sum_r e^{-U_r} \mathrm{diag}\,(c_1^{(r)}, c_2^{(r)}, c_3^{(r)}) \cdot \begin{pmatrix} 1 & 1 & 1 \\ 1 & 1 & 1 \\ 1 & 1 & 1 \end{pmatrix} \cdot \mathrm{diag}\,(c_1^{(r)}, c_2^{(r)}, c_3^{(r)}) \ . \tag{25}$$

- Each instanton makes one particular combination of ν_L massive. A hierachy among 3 eigenvalues naturally appears for 3 instantons with different $exp(-U_r)$

[a] L.I.,B.Schellekens, A. Uranga, hep-th/0704.1079

- One can get a structure consistent with experiment [b] . For example, if two instantons $D2^{(2)}$ and $D2^{(3)}$ dominate with $exp(-ReU_3)/exp(-ReU_2) \sim 5$ and

$$(c_1^{(2)}, c_2^{(2)}, c_3^{(2)}) = \frac{1}{\sqrt{3}}(1,1,1)\ ,(c_1^{(3)}, c_2^{(3)}, c_3^{(3)}) = \frac{1}{\sqrt{2}}(0,-1,1)\ . \tag{26}$$

- Asuming small mixing in the charged lepton mass matrix one obtains hierarchy of masses and (aproximately) Tri-bimaximal mixing for PMNS matrix

$$U_{\rm tri} = \begin{pmatrix} \sqrt{2/3} & 1/\sqrt{3} & 0 \\ -1/\sqrt{6} & 1/\sqrt{3} & -1/\sqrt{2} \\ -1/\sqrt{6} & 1/\sqrt{3} & 1/\sqrt{2} \end{pmatrix}\ . \tag{27}$$

- This is just an example. More generaly, there are ranges of $c_a^{(i)}$ consistent with the data.

[b] Antusch,L.I.,Macri Arxiv:0760.2132 [hep-ph].

Neutrino Mass Matrices: See-Saw

- Again, for $Sp(2)$ instantons ν_R-masses flavor dependence simplified:

$$M^{\nu_R}{}_{ab} \ = \ 2M_s \sum_r d_a^{(r)} d_b^{(r)} \ e^{-U_r} \tag{28}$$

- Again, each instanton makes one particular combination of ν_R massive. A hierarchy among 3 eigenvalues naturally appears for 3 instantons with different $exp(-U_r)$.

-

$$M^{\nu_L}(\text{see-saw}) \ = \ \frac{<\overline{H}>^2}{2M_s} h_D^T (\sum_r \ d_a^{(r)} d_b^{(r)} \ e^{-U_r})^{-1} \ h_D \tag{29}$$

- This case is more model dependent, structure depends on ordinary Yukawa coupling contant matrices h_D. Again, agreement with expermimant may be obtained e.g.for d_a aligned with diagonal h_D

- Both Weinberg and see-saw mechanisms may be simultaneously present

When does the mechanism work?

- i) The SM group should be extended by a $U(1)_{B-L}$ gauge symmetry. There are ν_R 's in the massless spectrum to begin with.
- ii) The gauge boson $U(1)_{B-L}$ gets a Stuckelberg mass combining with an axion-like closed string field.
- iii) String Dp-instantons exist with the appropriate fermionic zero modes.
- Local examples exist in which indeed the neutrino mass is generated [a] It requires D2-branes wrapping rigid cycles.
- This is stringy mechanism, because these are string instanton effects with no obvious field theory counterpart.
- The same mechanism exists in other known compactifications like Type IIB orientifolds and heterotic (with $U(1)$ bundles).

[a]Cvetic,Richter,Weigand hep-th/0703028.

Other Interesting Instanton-induced Operators

- It is clear that in general OTHER instantons may exist generating operators violating some massive $U(1)$ symmetry:

$$e^{-S_{Ins}} \; \phi....\phi \tag{30}$$

- Examples are
  - The μ-term in the MSSM

$$e^{-S_{Ins}} \; H\overline{H} \tag{31}$$

  - Some Yukawa couplings which may be forbidden perturbatively (e.g., for the 1-st generation). For example Lepton Yukawas in the LR-symmetric D3-brane model.
  - R-parity violating couplings in MSSM : $L\bar{H}$, UDD .
  - Superpotential couplings involving hidden sector fields, possibly usefull in fixing moduli and/or breaking SUSY[a]

[a] Florea et al.; Akerblom et al. 2006

VII- FLUXES AND SUSY BREAKING

SUSY breaking and fluxes

- Imagine we succed in building a string vacuum with the structure of the MSSM. Does string theory give us information about the structure of SUSY breaking soft terms?.

$$
\begin{aligned}
L_g &= \frac{1}{2}\sum_a M_a \,\lambda_a\lambda_a \,+\, h.c. \qquad (32)\\
L_{m^2} &= -m_{H_d}^2|H_d|^2 - m_{H_u}^2|H_u|^2 - m_{Q_{ij}}^2 Q_iQ_j^* - m_{U_{ij}}^2 U_iU_j^* \\
&- \; m_{D_{ij}}^2 D_iD_j^* - m_{L_{ij}}^2 L_iL_j^* - m_{E_{ij}}^2 E_iE_j^* \\
L_{A,B} &= -\,A_{ij}^U Q_iU_jH_u \,-\, A_{ij}^D Q_iD_jH_d \,-\, A_{ij}^L L_iE_jH_d \,-\, B\,H_dH_u \,+\, h.c.
\end{aligned}
$$

- There are plenty of possibilities for SUSY breaking and its mediation. In essentially all approaches there is a hidden sector of SM singlets in which SUSY breaking resides.
- In string theory the Kahler moduli T_i, dilaton S and complex structure U_i fields are natural candidates to form part of the hidden sector.
- One simple possibility is to asume that the source of SUSY breaking resides in the auxiliary fields of the dilaton/moduli fields , e.g. , F_S, F_{T_i} .
- If we have a knowledge of the Kahler potential and gauge kinetic function one can then derive predictions for the soft terms. In the Heterotic it was found:
  - $F_T \neq 0$ **: MODULUS DOMINANCE**

    This gives rise to a vanishing c.c. to leading order , which is a nice point. However no-soft terms appear (again to leading order).
  - $F_S \neq 0$ **: DILATON DOMINANCE**

This gives rise to interesting universal soft terms:

$$m^2 = \frac{1}{3}|M|^2 \; ; \; A^{ijk} = -h^{ijk} M_a \tag{33}$$

– However no obvious microscopic source for such $F_S \neq 0$ was found in the context of the heterotic string.

- As we have seen, Type II orientifolds, e.g. Type IIB orientifolds compactified on a CY offer new possibilities for the embedding of the SM in string theory.
- In a different developement it has been realized the important role played by antisymmetric field fluxes in Type IIB orientifold compactifications.
- Fluxes in Type IIB orientifold theories may fix both the dilaton and the complex-structure moduli M_i. Including non-perturbative effects depending on the volume moduli T_i all the moduli in these compactifications could possibly be determined.
- Here we are going to discuss another consequence of the presence of fluxes : they may GENERATE SUSY BREAKING SOFT TERMS.

- Note that if SUSY is found at LHC , specific relationships among soft parameters could be tested experimentally. Thus finding a theory of soft terms is a challenge we cannot refuse!!

Closed string fluxes in IIB

- Type II string theory contains antisymmetric fields from the RR and NS sectors. Their field strengths may be non-vanishing and the corresponding fluxes through closed surfaces are quantized.
- These provide for new (discrete) degrees of freedom in each compactification.
- The best understood case is that of 3-form fluxes in Type IIB orientifolds. There are NS H_3 and RR F_3. They would verify for any 3-cycle Σ in the CY

$$\frac{1}{(2\pi)^2\alpha'}\int_\Sigma F_3 \in \mathbf{Z} \; ; \; \frac{1}{(2\pi)^2\alpha'}\int_\Sigma H_3 \in \mathbf{Z} \qquad (34)$$

- These fluxes contribute positively to the vacuum energy. That lead to certain no-go theorems stating that no IIB compactifications with fluxes are consistent with equations of motion.

- Fluxes also contribute to tadpoles of RR C_4 form due to the CS coupling

$$\int_{M_4 \times CY} H_3 \wedge F_3 \wedge C_4 \tag{35}$$

One has in fact

$$N_{flux} = \frac{1}{(4\pi^2\alpha')^2} \int_{CY} H_3 \wedge F_3 = \frac{1}{(4\pi^2\alpha')^2} \frac{1}{(S+S^*)} \int_{CY} G_3 \wedge \overline{G_3} \tag{36}$$

where S is the complex dilaton and

$$G_3 = F_3 - iSH_3 \tag{37}$$

- Both things (vacuum energy and RR tadpoles) may be cancelled in the orientifold case, the orientifolds have negative tension and negative charge. Solutions of equations of motion are obtained (Giddings, Kachru Polchinski) if G_3 is imaginary self-dual (ISD):

$$*_6 G_3 \; = \; i \, G_3 \tag{38}$$

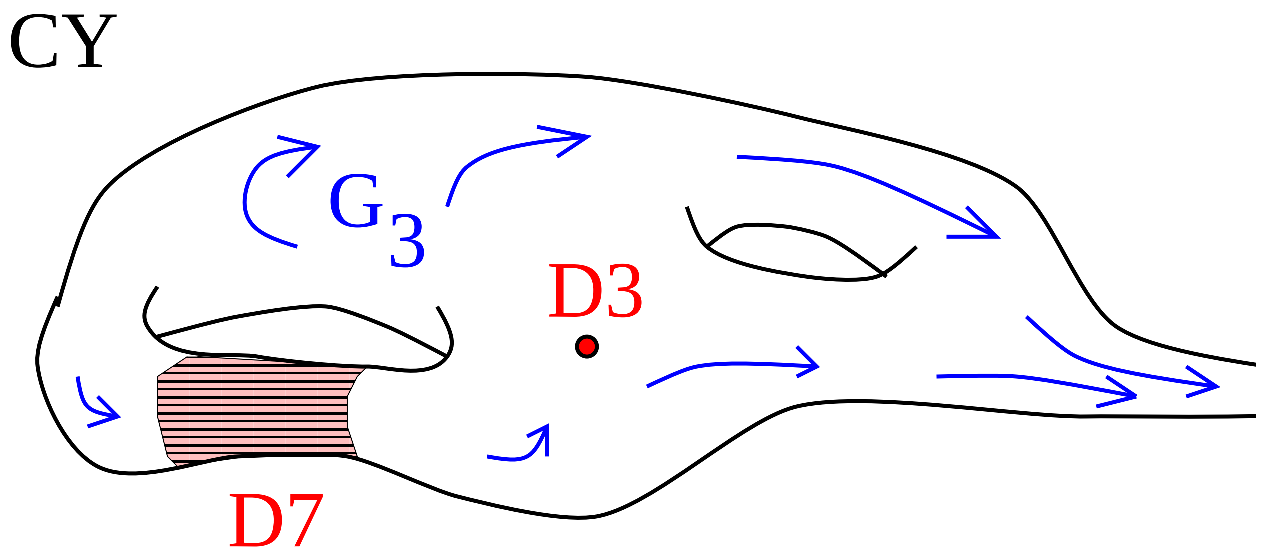
- The presence of fluxes generates a superpotential for the complex structure U_i and dilaton S fields

$$W_{GVW} \; = \; \int_{CY} G_3 \wedge \Omega^{CY} \tag{39}$$

where Ω^{CY} is the holomorphic 3-form of the CY, which contains dependence on complex structure fields, (e.g. for tori $\Omega = (dx_1 + \tau_1 dy_1) \wedge (dx_2 + \tau_2 dy_2) \wedge (dx_3 + \tau_3 dy_3)$).

- Minimization of the effective potential generically fixes the complex structure and dilaton fields U_i, S. (But not the Kahler moduli T_i).

D3/D7-branes and fluxes



- We consider $D3$-branes localized at points in CY.
- D7-branes wrap 4-cycles over the CY
- G_3 fluxes are present in the CY.

ISD and IASD 3-form fluxes

- The flux G_{mnp} transforms as a 20-dimensional reducible $SO(6)$ representation, decomposing as $\mathbf{20} = \overline{\mathbf{10}} + \mathbf{10}$ (imaginary self-dual (ISD) $G^{+}_{(3)}$ and imaginary anti self-dual (IASD) $G^{-}_{(3)}$ parts), respectively

$$G^{\pm}_{(3)} \;=\; \frac{1}{2}(G_{(3)} \mp i *_6 G_{(3)}) \;\; ; \;\; *_6 G^{\pm}_{(3)} \;=\; \pm i G^{\pm}_{(3)} \tag{40}$$

- It is useful to classify the components of the ISD and IASD parts of G_3 according to their behavior under $SU(3)$ (Graña, Polchinski), $\mathbf{10} = \mathbf{6} + \mathbf{3} + \mathbf{1}$

	ISD			IASD	
$SU(3)$ rep.	Form	Tensor	$SU(3)$ rep.	Form	Tensor
$\overline{1}$	$(0,3)$	$G_{\bar{1}\bar{2}\bar{3}}$	1	$(3,0)$	G_{123}
$\overline{6}$	$(2,1)_P$	$S_{\bar{i}\bar{j}}$	6	$(1,2)_P$	S_{ij}
$\overline{3}$	$(1,2)_{NP}$	A_{ij}	3	$(2,1)_{NP}$	$A_{\bar{i}\bar{j}}$

- Here one defines (and similarly for $S_{\bar{i}\bar{j}}, A_{ij}$)

$$S_{ij} = \frac{1}{2}(\epsilon_{ikl}G_{j\bar{k}\bar{l}} + \epsilon_{jkl}G_{i\bar{k}\bar{l}}) \ ; \ A_{\bar{i}\bar{j}} = \frac{1}{2}(\epsilon_{\bar{i}\bar{k}\bar{l}}G_{kl\bar{j}} - \epsilon_{\bar{j}\bar{k}\bar{l}}G_{kl\bar{i}}) \quad (41)$$

- Subindex P=primitive $\leftrightarrow G_3 \wedge J = 0$. (In general Non-primitive fluxes not present in CY, since if $G_3 \wedge J \neq 0$ 5-cycles should exist, and they do not in CY).
- (0,3) ISD ((3,0) IASD) flux contributes to RR-charge with same sign as $D3$ ($\overline{D3}$)-branes
- (2,1) fluxes preserve SUSY whereas (0,3) do not.

Effect of fluxes on D3-brane fields

- Let us consider a stack of $D3$ branes at a (smooth) point in the CY (the results for an orbifold singularity may be obtained by projection). In terms of $N = 1$ SUSY we have super-Yang-Mills and 3 chiral adjoint multiplets Φ_i, i=1,2,3.

- The vev of the 3 worldvolume scalars ϕ_i correspond to the three complex transverse coordinates $x^m = 2\pi\alpha' \phi^m$

- The effective action in the presence of fluxes may be obtained by expanding the Dirac-Born-Infled action (as extended by Myers to non-Abelian case):

$$S_{BI} = -\mu_3 \int d^4x \, \mathrm{Tr} \left(e^{-\phi} \sqrt{- \det \left(P \left[E_{\mu\nu} + E_{\mu m} (Q^{-1} - \delta)^{mn} \, E_{n\nu} \right] + \sigma \, F_{\mu\nu} \right) \, det(Q)} \right.$$

$$\begin{aligned} E_{MN} &= G_{MN} - B_{MN} \\ Q^m{}_n &= \delta^m{}_n + i\sigma \left[\phi^m, \phi^p \right] E_{pn} \\ \sigma &= 2\pi\alpha' \end{aligned} \tag{43}$$

as well as the Chern-Simons action S_{CS} which includes RR couplings (plus fermionic terms).

- Plugging the closed string background one obtains [a] the SUSY breaking soft terms (we take $A_{mn} = 0$).

$$
\begin{aligned}
m_{ij}^2 &= \frac{g_s}{6}(|G_{123}|^2 + \sum_{ij}|S_{ij}|^2 - Re(G_{123}G_{\bar{1}\bar{2}\bar{3}} + \frac{1}{4}S_{lk}S_{\overline{lk}})) \\
A^{ijk} &= -h^{ijk}\frac{g_s^{1/2}}{\sqrt{2}}G_{123} \\
M^a &= \frac{g_s^{1/2}}{\sqrt{2}}G_{123} \\
\mu_{ij} &= -\frac{g_s^{1/2}}{2\sqrt{2}}S_{ij}
\end{aligned}
\tag{44}
$$

- Note that if only ISD fluxes ($G_{\bar{1}\bar{2}\bar{3}}$, $S_{\overline{lk}}$) are present, no soft terms are

[a] M. Graña hep-th/0209200, P. Camara et al.hep-th/0311241; Graña et al. hep-th/0312232.

generated!. This is like in the heterotic modulus dominance situation.

- On the other hand in the presence of IASD $(3,0)$ fluxes G_{123} one gets

$$m^2 = \frac{g_s}{6}|G_{123}|^2 \; ; \; M^a = \frac{g_s^{1/2}}{\sqrt{2}} G_{123} \; ; \; A^{ijk} = -h^{ijk} \frac{g_s^{1/2}}{\sqrt{2}} G_{123} \quad (45)$$

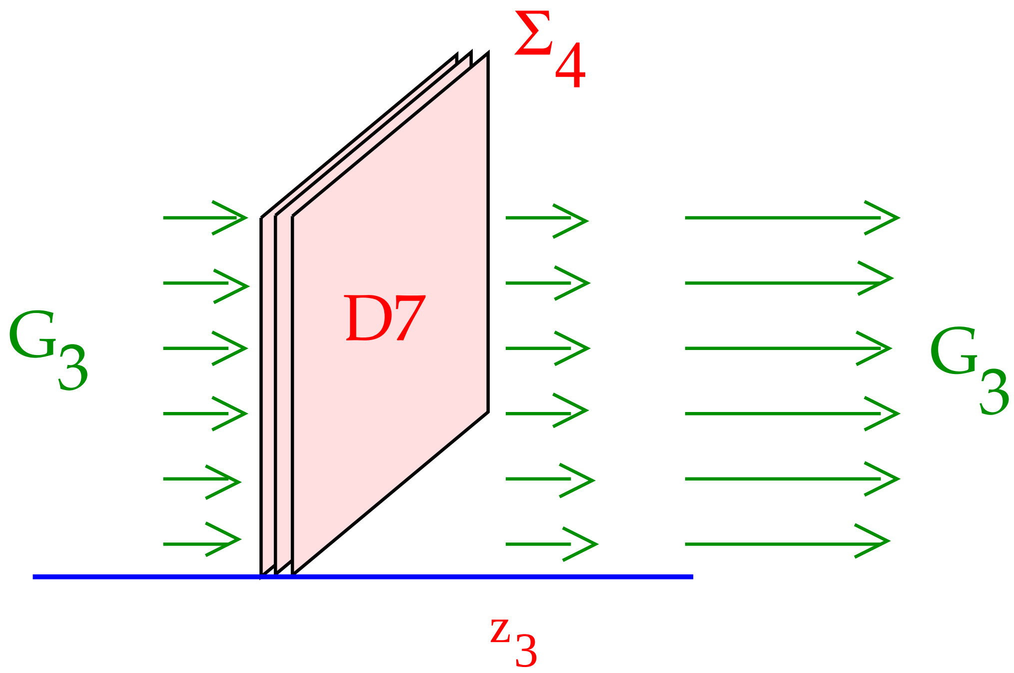
- Note the relationships

$$A^{ijk} = -h^{ijk} M_a \; ; \; m^2 = \frac{1}{3}|M|^2 \quad (46)$$

- They correspond to dilaton dominated SUSY-breaking previously discussed in heterotic context.
- Note that anti-D3branes in $(0,3)$ backgrounds have these soft terms.
- On the other hand only ISD fluxes are known to solve the equations of motion.
- One may conclude that, if one wants to have SUSY breaking induced by ISD fluxes, locating the MSSM at at $D3$ branes is not a good idea. Let us locate the SM at $D7$-branes.

Soft terms on D7-branes

- In this case the D7-branes wrap 4-cycles Σ_4 in the CY. The local tangent symmetry is $SO(4) \times SO(2)$ instead of $SO(6)$.



- To simplify matters one restricts the study to local geometries $T^4 \times \mathbf{C}$ or

$K3 \times \mathbf{C}$ and assume constant backgrounds over Σ_4. Similar features expected in more complicated geometries.

- We will asume our $D7$ brane transverse coordinate is x_3. There are now two types of chiral matter fields when reducing to 4 dimensions:
  - ϕ_3 which parametrizes the location of the $D7$-branes in transverse space
  - $\phi_{1,2}$ which come from the zero modes of the gauge fields in 8 dimensions.
- Doing a similar expansion of the $S_{DBI} + S_{CS}$ as in the $D3$ case one finds soft terms depending on the fluxes as follows [a]:

$$
\begin{aligned}
m^2_{1\bar{1}} &= m^2_{2\bar{2}} = 0 \; ; \; B_{ij} \; = \; 0 \; , \; i,j \neq 3 \\
m^2_{3\bar{3}} &= \frac{g_s}{18}\left(|G_{\bar{1}\bar{2}\bar{3}}|^2 + \frac{1}{4}|S_{\bar{3}\bar{3}}|^2 + \frac{1}{4}\sum_{i,j=1,2} |S_{ij}|^2 \right) \\
B_{33} &= \frac{g_s}{9}\left(\frac{1}{4}(S_{12})^{*2} - \frac{1}{2}(G_{\bar{1}\bar{2}\bar{3}})^*(S_{\bar{3}\bar{3}})^* - \frac{1}{4}(S_{22})^*(S_{11})^* \right)
\end{aligned}
$$

[a] P. Camara et al hep-th/0408036; Lust et al. hep-th/0406092.

$$
\begin{aligned}
A^{ijk} &= -h^{ijk}\,\frac{g_s^{1/2}}{3\sqrt{2}}\,(G_{\bar{1}\bar{2}\bar{3}})^* \\
M^a &= \frac{g_s^{1/2}}{3\sqrt{2}}\,(G_{\bar{1}\bar{2}\bar{3}})^* \\
\mu_{33} &= -\frac{g_s^{1/2}}{6\sqrt{2}}(S_{\bar{3}\bar{3}})^* \\
\mu_{ij} &= -\frac{g_s^{1/2}}{6\sqrt{2}}S_{ij}\ ,\ i,j=1,2
\end{aligned}
\qquad (47)
$$

The case of ISD fluxes

- These include fluxes $(0,3)$ $(G_{\bar{1}\bar{2}\bar{3}})$ and $(2,1)$ $(S_{\bar{i}\bar{j}})$. One finds

$$m^2_{\Phi^3_{77}} \;=\; \frac{g_s}{18}\,|G_{\bar{1}\bar{2}\bar{3}}|^2 \;;\; M^{(77)} \;=\; \frac{g_s^{1/2}}{3\sqrt{2}}\,(G_{\bar{1}\bar{2}\bar{3}})^* \;;\; A^{ijk\,(77)} \;=\; -h^{ijk}\,\frac{g_s^{1/2}}{3\sqrt{2}}\,(G_{\bar{1}\bar{2}\bar{3}})^* \tag{48}$$

$$\mu_{(77)} \;=\; -\frac{g_s^{1/2}}{6\sqrt{2}}(S_{\bar{3}\bar{3}})^* \;;\; B_{33} \;=\; -\frac{g_s}{18}\,(G_{\bar{1}\bar{2}\bar{3}})^*(S_{\bar{3}\bar{3}})^* \;=\; 2M\mu_{(77)} \tag{49}$$

- Interestingly enough, for ISD backgrounds the scalar potential is positive definite

$$V_{ISD} \;=\; |-M_{77}{}^*\Phi^3_{77}{}^* - \mu_{(77)}\Phi^3_{77} + \Phi^1_{77}\Phi^2_{77} + \Phi_{73}\Phi_{37}\,|^2 \;=\; |-M_{77}{}^*\Phi^3_{77}{}^* + \partial_{\phi^3}W\,|^2 \tag{50}$$

- Note that the scheme is very constrained yielding relationships

$$A^{ijk\,(77)} \;=\; -h^{ijk}M^{(77)} \;;\; m^2_{\Phi^3_{77}} \;=\; |M^{(77)}|^2 \;;\; B_{33} \;=\; 2M\mu_{(77)} \tag{51}$$

Properties of soft terms on D3- and D7-branes

- On D3-branes ISD fluxes do not give rise to soft terms.
- On D7-branes ISD fluxes $(0,3)$ give rise to soft terms. The same is expected for fields at intersecting D7-branes.
- The fact that ISD fluxes give rise to soft terms on D7-branes is important since such fluxes verify eqs. of motion in CY compactification (GKP).
- Suggest to locate MSSM fields at D7-branes (or their intersections).
- Note that μ-terms are generated of order of soft masses. That might solve the μ-problem in MSSM-like models.
- Fluxes may be added to the 3-generation $Z_2 \times Z_2$ orientifold model cancelling global tadpoles[a].

[a] Marchesano, Shiu hep-th/0409132.

Understanding soft terms from the effective action

- It turns out that the switching on fluxes correspond to non-vanishing auxiliary fields for dilaton S /moduli T_i.
- Consider a simplified case of a $N=1$ CY compactification with a single (large) Kahler modulus T

$$K \;=\; -\log(S+S^*) \;-\; 3\log(T+T^*)\;;\; W \;=\; \int G_{(3)}\wedge\Omega \quad (52)$$

- Here W is the superpotential generated by the flux. The supergravity auxiliary field for a chiral field with metric $K_{i\bar{j}}$ is given by

$$\overline{F}^{\bar{i}} \;=\; \exp(K/2M_p^2)\; K^{\bar{i}j}\; D_jW/M_p^2 \quad (53)$$

- Applying to the above Kahler potential and superpotential one obtains

$$F^S \;=\; (S+S^*)^{1/2}\,(T+T^*)^{-3/2}\int G_{(3,0)}\wedge\Omega$$

$$F^T \quad = -(S+S^*)^{-1/2}\,(T+T^*)^{-1/2} \int G_{(0,3)} \wedge \Omega \qquad (54)$$

- This shows that a $(3,0)$ flux corresponds to a non-vanishing auxiliary field for the complex dilaton S, whereas a $(0,3)$ corresponds to the overall Kähler field T. One then has

$$m_{3/2}^2 = e^K |W|^2 = (S+S^*)^{-1}\,(T+T^*)^{-3}\,|\int G_{(0,3)} \wedge \Omega\,|^2 \,. \qquad (55)$$

- One can now consider the presence of D7 branes in the CY . Consider the case of a toroidal/orbifold orientifold compactification with $D7^i$-branes wrapping the T^4 transverse to the $i-th$ complex plane. One has gauge kinetic functions

$$f_{7_i} \;=\; T_i \qquad (56)$$

- We saw that the two types of matter fields in the worldvolume of D7-branes turn out to have Kahler metrics (ignoring here c.s. fields):

$$K_{ii}^{i} = \frac{1}{S+S^*} \; ; \; K_{ii}^{j} = \frac{1}{T_k+T_k^*} \quad , \quad i \neq j \neq k \tag{57}$$

- On the other hand for states for intersecting $D7_i - D7_j$-branes we already mentioned (for the T-dual of intersecting D6-branes):

$$K_{ij} = \frac{1}{(S+S^*)^{1/2}(T+T^*)^{1/2}} \tag{58}$$

- In terms of the dependence on the overall Kahler modulus $T = T_i$ one can summarize for the gauge kinetic function and the matter metrics the result

$$f = T \; ; \; K_{\xi} = \frac{1}{s^{1-\xi} t^{\xi}} \tag{59}$$

where $t = T + T^*$, $s = S + S^*$ and the 'modular weights' $\xi = 0, 1/2, 1$ depending the origin of the matter field.

- One can argue that for large volume these 3 types of matter fields coming from D7-branes are general, not particular to the toroidal case.

- The $D7$ branes generically may have magnetic flux F_i in their worldvolume in order to obtain chirality. As we saw this flux modifies the Kahler moduli in a toroidal setting. One finds for $D7^i$ branes within the same stack the result

$$K_{(7^i7^i)_j} \; = \; \frac{1}{t^k} \, |\frac{1 + iF^k}{1 + iF^j}| \;\; ; \;\; K_{(7^i7^i)_i} \; = \; \frac{1}{s} \, (1 + |F^j F^k|) \, , \tag{60}$$

where $i \neq j \neq k$ label the 3 2-tori and F^i is the magnetic flux going through the i-th 2-torus which may be written as

$$F^i \; = \; n^i (\frac{st_i}{t_j t_k})^{1/2} \, , \tag{61}$$

with n^i quantized integer fluxes.

- For states coming from open strings in between (magnetized) branes $D7^a$, $D7^b$ wrapping different 4-tori we already mentioned (for the T-dual case) one has

$$K_{7^a7^b} \; = \; \frac{1}{(st_1t_2t_3)^{1/4}} \Pi_{j=1}^3 \, u_j^{-\theta^j_{ab}} \sqrt{\frac{\Gamma(\theta^j_{ab})}{\Gamma(1-\theta^j_{ab})}} \, , \qquad (62)$$

where u_j are the real parts of the complex structure moduli, Γ is the Euler Gamma function and

$$\theta^j_{ab} \; = \; \arctan(F^j_b) - \arctan(F^j_a) \, . \qquad (63)$$

- The gauge kinetic functions are also modified in the presence of magnetic fluxes as

$$Ref_i^a \; = \; T_a \, (1 \, + \, |F^j_a F^k_a|) \, . \qquad (64)$$

- These results apply for a toroidal and/or orbifold setting. However we can try to model out what could be the effect of fluxes in a more general setting following the above structure. To model out the possible effect of fluxes we consider the limit with $t_i = t$ and diluted fluxes $|F_i| = F$, i.e. large t and ignore the dependence on the complex structure u_i fields. Then one gets from the above formulae:

$$K^i_{(7^i7^i)_j} = \frac{1}{t} \; ; \; K_{(7^i7^i)_i} = \frac{1}{s}(1 + a_i\frac{s}{t}) \, , \tag{65}$$

$$K_{7^a7^b} = \frac{1}{(s^{1/2}t^{1/2})}(1 + c_{ab}\frac{s^{1/2}}{t^{1/2}}) \, , \tag{66}$$

$$\mathrm{Re} f_i = t + a_i s \, , \tag{67}$$

where a_i, c_{ab} are constants (including the flux quanta) of order one.

- These three formulae may again be summarised by:

$$K_{matter} = \frac{1}{s^{(1-\xi)}t^{\xi}} \times (1 + c_\xi (s/t)^{1-\xi}) = \frac{1}{s^{(1-\xi)}t^{\xi}} + \frac{c_\xi}{t} \, , \tag{68}$$

with c_ξ some flux-dependent constant coefficient whose value will depend on the modular weight ξ and the magnetic quanta.

- In the dilute flux $t \rightarrow \infty$ one recovers the fluxless case.

MSSM soft terms computation from modulus dominance

- Asume that we find a MSSM string vacuum with the SM particles iving on $D7$-branes.

- As we have shown this has the advantage that ISD fluxes breaking SUSY lead to SUSY-breaking soft terms already at leading order. This correspond to non-vanishing vevs for the auxiliary fields of Kahler moduli: modulus dominance.

- In the large Kahler moduli limit we can then compute the SUSY breaking soft terms under a number of simplifying assumptions:
  - Assume that eventually all moduli are stabilized with a vacuum energy close to zero.
  - Assume MSSM particle content and standard gauge coupling unification.
  - Insist that the D-brane configuration is such that there is at least one Yukawa coupling (top quark) of order g (the gauge coupling constant).

- Gauge coupling unification may be natural if the MSSM resides at F-theory 7-branes (rather than D7-branes) (Beasley,Heckman,Vafa; Donagie,Wijnholt).
- A further simplification is that we will consider a single local Kahler modulus t coupling to the MSSM brane system. t is not in general an overall volume modulus but a local modulus.
- Then, using the Kahler metrics and gauge kinetic function discussed above one can compute the MSSM SUSY-breaking soft terms in the standard way.

Soft term computation from the effective action

- Using the effective $N = 1$ supergravity Lagrangian one can compute soft terms as

$$
\begin{aligned}
M_i &= \frac{1}{2\mathrm{Re}\, f_i} F^M \partial_M f_i \ , \\
m_I^2 &= m_{3/2}^2 - \sum_{M,N} \bar{F}^{\bar{M}} F^N \partial_{\bar{M}} \partial_N \log(\tilde{K}_{I\bar{I}}) \ , \\
A_{IJL} &= F^M [\hat{K}_M + \partial_M \log(Y_{IJL}) - \partial_M \log(\tilde{K}_{I\bar{I}} \tilde{K}_{J\bar{J}} \tilde{K}_{L\bar{L}})] \\
B &= \left(F^m \left[\hat{K}_m + \partial_m \log \mu \ - \ \partial_m \log(K_{H_u} K_{H_d}) \right] - m_{3/2} \right)
\end{aligned}
$$

Here F^M are auxiliary fields of moduli, $\tilde{K}_{I\bar{I}}, \hat{K}_M$ the metric of matter and moduli and f_i gauge kinetic functions.

- In Type IIB orientifolds the holomorphic perturbative superpotential is independent of the Kahler moduli so that the derivatives of $\widehat{Y}^{(0)}$ in the expression for A vanish.

- Using these formulae one then obtains [a] general soft terms as follows (the gaugino mass M is a free parameter)::

$$\begin{aligned} m_\alpha^2 &= (1-\xi_\alpha)|M|^2 \ , \quad \alpha = Q, U, D, L, E, H_u, H_d \ , \\ A_U &= -M(3-\xi_{H_u}-\xi_Q-\xi_U) \ , \\ A_D &= -M(3-\xi_{H_d}-\xi_Q-\xi_D) \ , \\ A_L &= -M(3-\xi_{H_d}-\xi_L-\xi_E) \ , \\ B &= -M(2-\xi_{H_u}-\xi_{H_d}) \ . \end{aligned} \tag{69}$$

where we have neglected for the moment the corrections from magnetic fluxes (i.e. dilute flux limit).

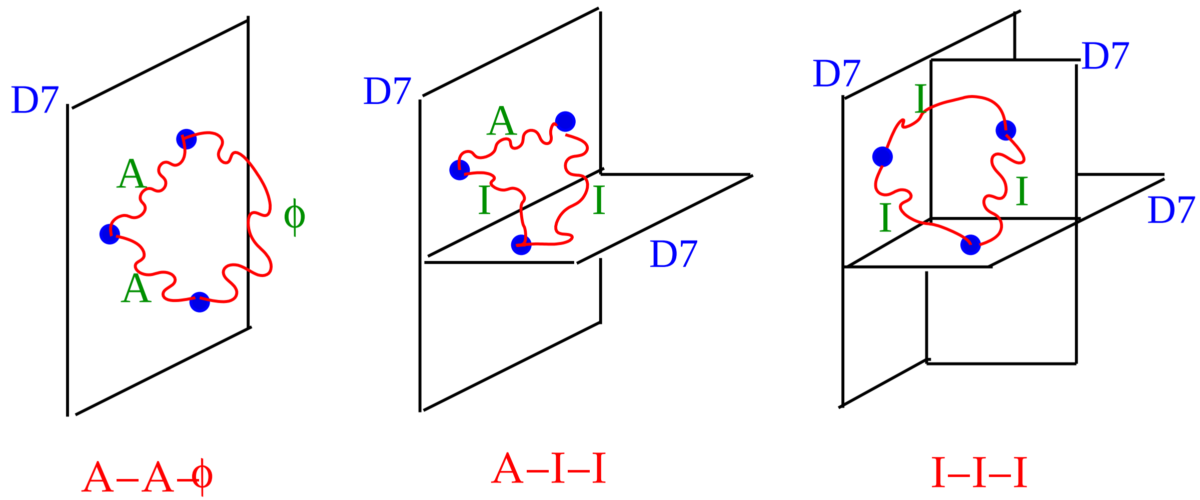
- Within the philosophy of gauge coupling unification one can assume unified modular weights:

$$\xi_f \ = \ \xi_Q = \xi_U = \xi_D = \xi_L = \xi_E \ . \tag{70}$$

[a] L. Aparicio,L.E.I., D. G.Cerdeño, hep-ph/0805.2943.

- On the other hand the Higgs fields could have different modular weight than fermion fields. So we will take $\xi_H = \xi_{H_u} = \xi_{H_d} = 0, 1, 1/2$.
- We have three type of 7-brane matter fields ϕ, A, I corresponding to modular weights 0,1,1/2 respectively.
- It turns out that there are only renormalizable couplings of three types.

$$(A - A - \phi) \ ; \ (I - I - A) \ ; \ (I, I, I) \qquad (71)$$



- They correspond to modular weights (1,1,0)), (1,1/2,1/2) and (1/2,1/2,1/2).

- The same types of Yukawa couplings exist in F-theory compactifications (Vafa et al.).
- For each of these three configurations the results for soft terms are shown in the table.

(ξ_L, ξ_R, ξ_H)	Coupling	M	m_L^2	m_R^2	m_H^2	A	B
$(1,1,0)$	(A-A-ϕ)	M	0	0	$\|M\|^2$	$-M$	$-2M$
$(1/2,1/2,1)$	(I-I-A)	M	$\frac{\|M\|^2}{2}$	$\frac{\|M\|^2}{2}$	0	$-M$	0
$(1/2,1/2,1/2)$	(I-I-I)	M	$\frac{\|M\|^2}{2}$	$\frac{\|M\|^2}{2}$	$\frac{\|M\|^2}{2}$	-3/2M	$-M$

Table 1: Modulus dominated soft terms for choices of modular weights ξ_α which are consistent with the existence of trilinear Yukawa couplings in 7-brane systems.

- Note that in the scenarios with couplings (A-A-ϕ) and (I-I-A) it is natural to assume that the Higgs field is identified with fields of type ϕ and A respectively and these are the cases displayed in the table.

- Concerning the B- parameter it is obtained assuming an explicit μ-term.
- The modular weights in the first line reproduces the results obtained from $(0,3)$ ISD fluxes. (No prediction from fluxes exists for particles at intersecting $D7$-branes).
- The modular weights in the third case are universal, a particular case of CMSSM boundary conditions.
- One can also estimate the possible effect of magnetic fluxes in the dilute limit $t \longrightarrow \infty$. Using the corrected formulae for gauge kinetic function and Kahler metrics discussed before one finds

Coupling	m_f^2	m_H^2	A	B
(A-A-ϕ)	0	$\vert M\vert^2(1-2\rho)$	$-M(1-\rho)$	$-2M(1-\rho)$
(I-I-A)	$\frac{\vert M\vert^2}{2}(1-\frac{3}{2}\rho_f)$	0	$-M(1-\rho_f)$	0
(I-I-I)	$\frac{\vert M\vert^2}{2}(1-\frac{3}{2}\rho_f)$	$\frac{\vert M\vert^2}{2}(1-\frac{3}{2}\rho_H)$	$-\frac{1}{2}M(3-\rho_H-2\rho_f)$	$-M(1-\rho_H)$

where

$$\rho = \frac{(c_H - as)}{t} \; ; \; \sigma = \frac{as}{t} \; ; \; \rho_f = \frac{c_f}{t^{1/2}} \; ; \; \rho_H = \frac{c_H}{t^{1/2}} \; , \tag{72}$$

EW symmetry breaking and SUSY spectrum

- One can take the above values for soft terms as boundary conditions at the GUT/String scale.

- The scheme is very predictive, there are only two free parameters M, μ . Once one imposses REW symmetry breaking one has just one free parameter M which sets the scale.

- One can solve (numerically) the renormalization group equations from the String to the Weak scale and compute the low energy SUSY spectrum and Higgs potential. (use SPheno2.2.3 and micrOMEGAs).

- In order to obtain apropriate $SU(2)_L \times U(1)_Y$ breaking one has to impose the conditions at the EW scale:

$$\mu^2 = \frac{-m_{H_u}^2 \tan^2\beta + m_{H_d}^2}{\tan^2\beta - 1} - \frac{1}{2} M_Z^2, \tag{73}$$

$$\mu B = \frac{1}{2}\sin 2\beta\,(m^2_{H_d} + m^2_{H_u} + 2\mu^2)\,, \tag{74}$$

where

$$tan\beta \; = < H_u/H_d > \quad . \tag{75}$$

- In addition there are a number of experimental constraints
  - LEP limits on SUSY particles and lightest Higgs boson.
  - $2.85 \times 10^{-4} \leq \; \mathrm{BR}(b \rightarrow s\gamma) \leq 4.25 \times 10^{-4}$ (Heavy Flavour Averaging Group).
  - BR$(B_s^0 \rightarrow \mu^+\mu^-) < 5.8 \times 10^{-8}$ at 95% c.l. (CDF)
  - Anomalous magnetic moment of the muon, $11.6 \times 10^{-10} \leq \; a_\mu^{\mathrm{SUSY}} \leq 43.6 \times 10^{-10}$.
  - WMAP limits on cold dark matter (applied to neutralino LSP), $0.1037 \leq \Omega h^2 \leq 0.1161$.

The intersecting (I-I-A) and (I-I-I) schemes

- These two cases may be analysed together by taking the Higgs modular weight ξ_H as a free parameter with $\xi_H = 1, 1/2$ reproducing the two limits . Then soft erms are

$$
\begin{aligned}
{m_{L,E,Q,U,D}}^2 &= |M|^2/2 \, , \\
{m_{H_u,H_d}}^2 &= (1-\xi_H)|M|^2 \, , \\
A_{U,D,L} &= -M(2-\xi_H) \, , \\
B &= -2M(1-\xi_H) \, .
\end{aligned}
\tag{76}
$$

- A sample of the resulting SUSY spectrum is shown in the figure as a function of $tan\beta$ for M=400 GeV and $\mu < 0$ for $\xi_H = 1/2, 1$.

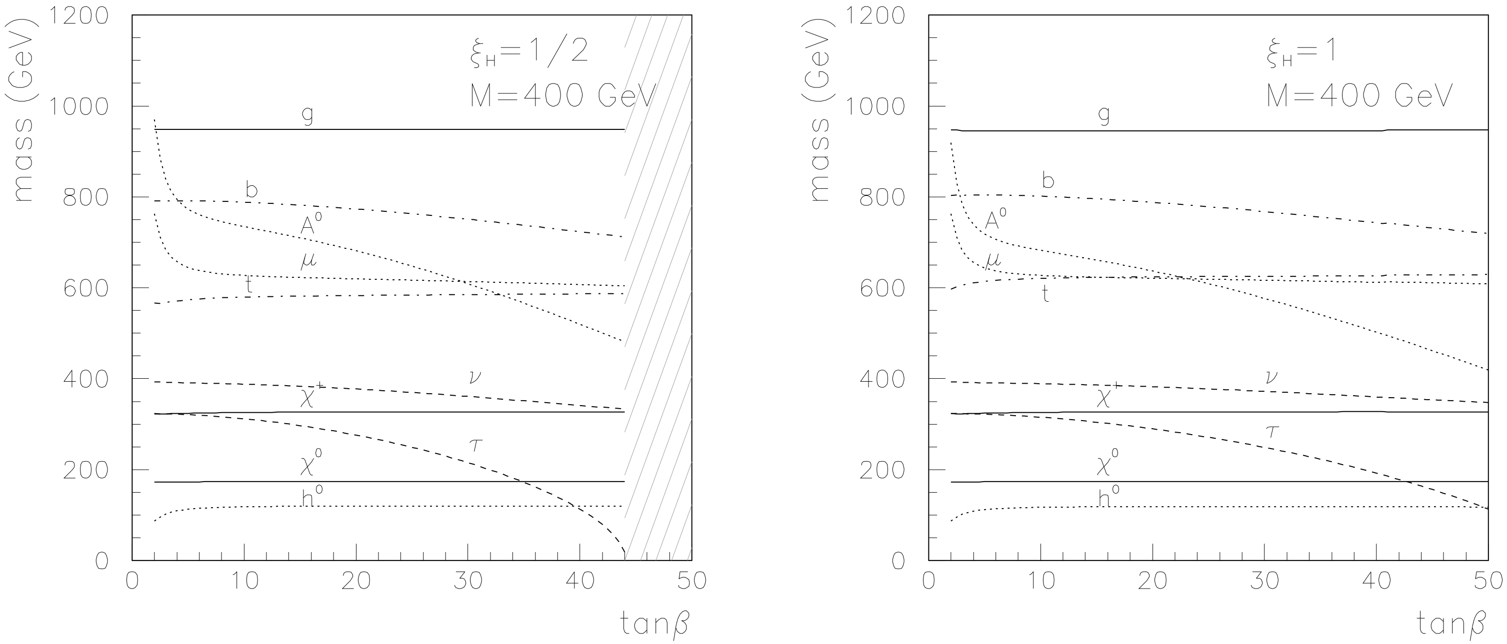


Figure 1: Low-energy supersymmetric spectrum as a function of $\tan\beta$ for $\xi_H = 1/2$, (left) and $\xi_H = 1$ (right) with $M = 400$ GeV and $\mu < 0$. The ruled area for large $\tan\beta$ is excluded by the occurrence of tachyons in the slepton sector.

- Note for not too large $tan\beta$ the lightest neutralino is the LSP. For large tanβ the lightest stau becomes lighter (and even tachyonic).

- This is because the Yukawa coupling goes like $1/cos\beta$ and r.g.e. decreases the mass^2 for large Yukawa. One thus has tan$\beta \leq 45$ GeV (for $\xi_H = 1/2$) and tan$\beta \leq 55$ (for $\xi_H = 1$).
- The effect of various experimental constraints is shown in the next figure for $\xi_H = 0.5, 0.6, 0.8, 1$ (condition for consistent radiative EW breaking given by a line).

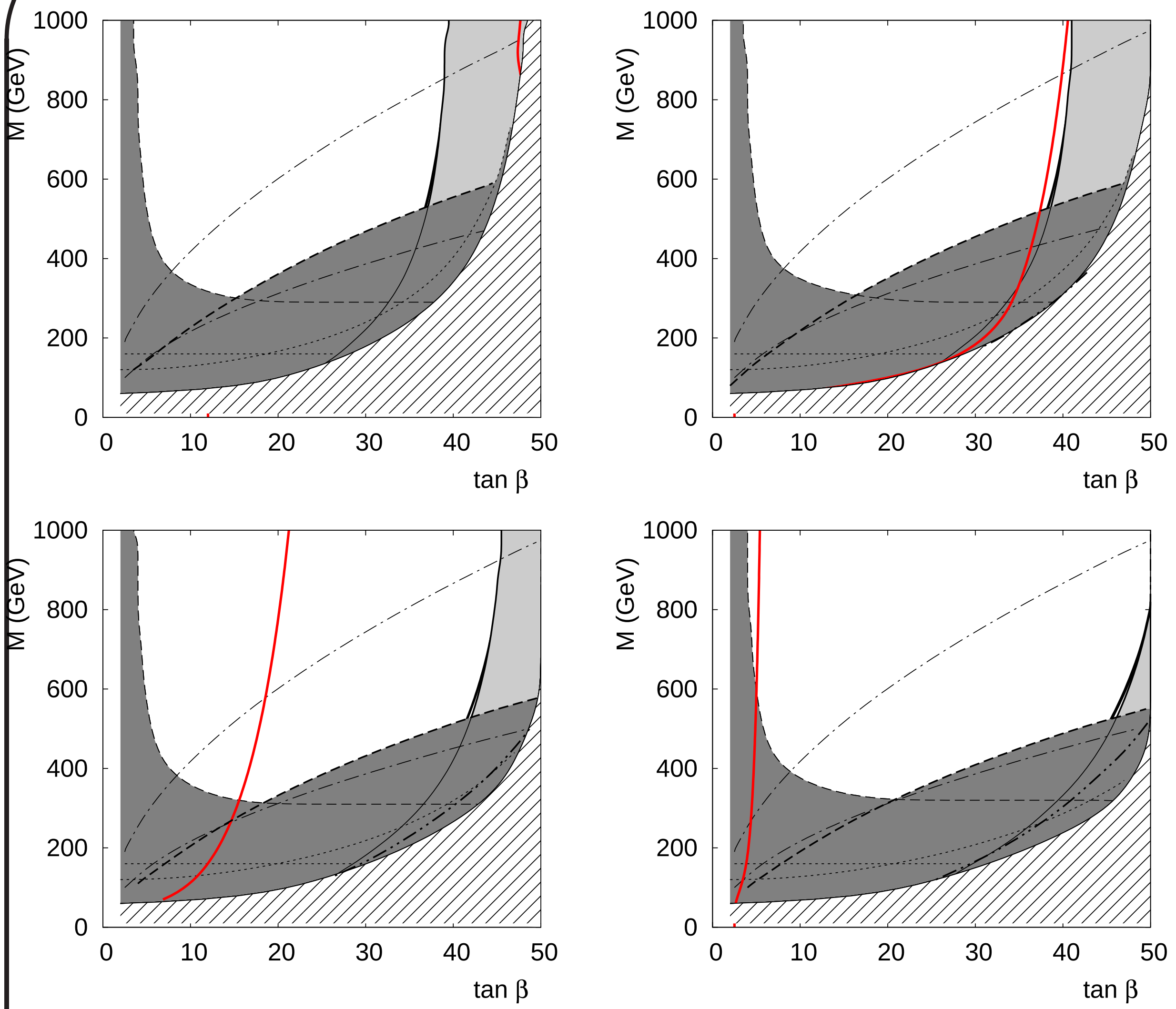


Figure 2: Effect of the various experimental constraints on the $(M, \tan\beta)$ plane for cases with $\xi_H = 0.5, 0.6, 0.8,$ and 1, from left to right and top to bottom. Dark grey regions correspond to those excluded by any experimental bound. Also shown is CONSISTENCY WITH REW BREAKING LINE.

- In order to get neutralino dark matter in agreement with WMAP results one should be in the coannihilation region with $m_{\chi^0} \simeq m_{\tilde{\tau}}$.

- On the other hand in order to achieve correct EW symmetry breaking in this coannihilation region one needs $\xi_H \simeq 0.6$ so that

$$\begin{aligned} {m_f}^2 &= 1/2\,|M|^2\,, \\ {m_H}^2 &\approx (1/2\,-\,0.1)\,|M|^2\,, \\ A_{U,D,L} &\approx (-3/2\,+\,0.1)\,M\,, \\ B &\approx (-1\,+\,0.2)\,M\,. \end{aligned} \tag{77}$$

- This in very close to the configuration with all particles residing at intersecting 7-branes. The small deviations may be atributed to subleading corrections .

- In particular those coming from the presence of magnetic fluxes . Small flux corrections with $\rho_H \simeq 0.17$, $\rho_f \simeq 0$ may account for these deviations.

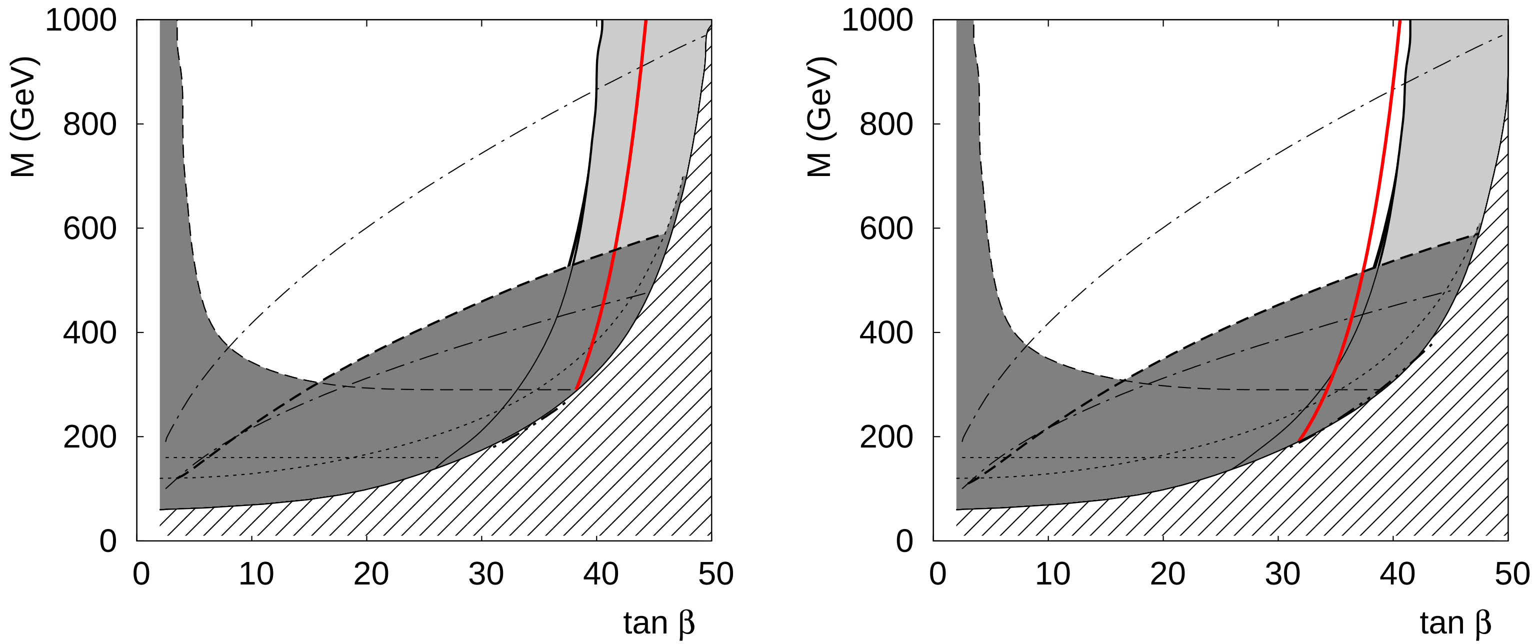
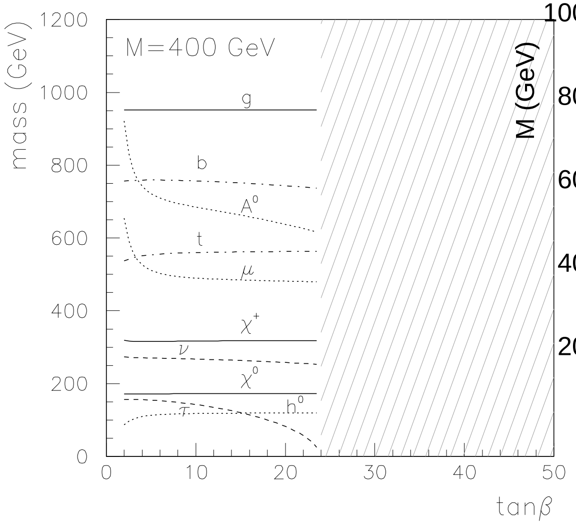


Figure 3: Flux corrections for $\rho_f = 0$ and $\rho_H = 0.1$ (left) and $\rho_H = 0.2$ (right).

The bulk 7-branes (A-A-ϕ) case

- The other configuration with Yukawa couplings with structure (A-A-ϕ) one finds spectra and experimental constraints very restricted:



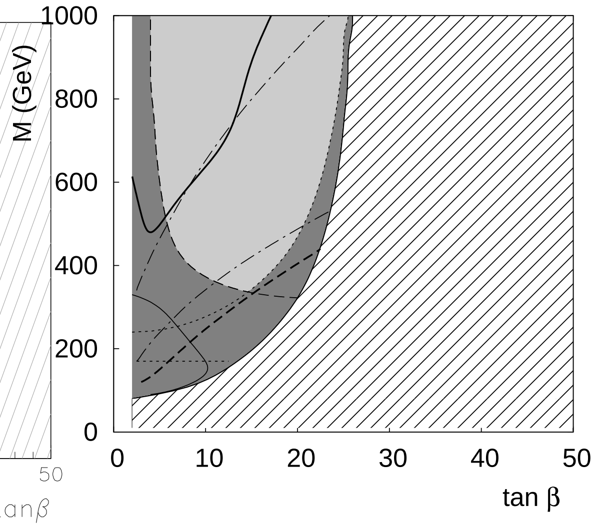


Figure 4: Left) Low-energy supersymmetric spectrum as a function of $\tan\beta$ for case (A-A-ϕ). Right) Effect of the various experimental constraints on the $(M, \tan\beta)$ plane for case (A-A-ϕ).

- The LSP is the stau. One does not get consistent REW.

LHC

- Making use of the missing energy signal for squarks and gluinos LHC will be able to test the intersecting brane scheme for

Int. Lumin.	M	$m_{\tilde{q}}$	$m_{\tilde{g}}$	$m_{\chi^0} \simeq m_{\tilde{\tau}}$
1 fb^{-1}	≤ 650	≤ 1.3	≤ 1.5	≤ 300
10 fb^{-1}	≤ 900	≤ 1.8	≤ 2.0	≤ 400

Spectra: