

# Soft Supersymmetry Breaking Terms in String Compactifications

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This talk is based on

[hep-th/0609180](#) (JC, D.Cremades, F. Quevedo)

[hep-th/0610129](#) (JC, S. Abdussalam, F. Quevedo, K. Suruliz)

[hep-ph/061xxxx](#) (JC, D. Cremades)

and also uses results from

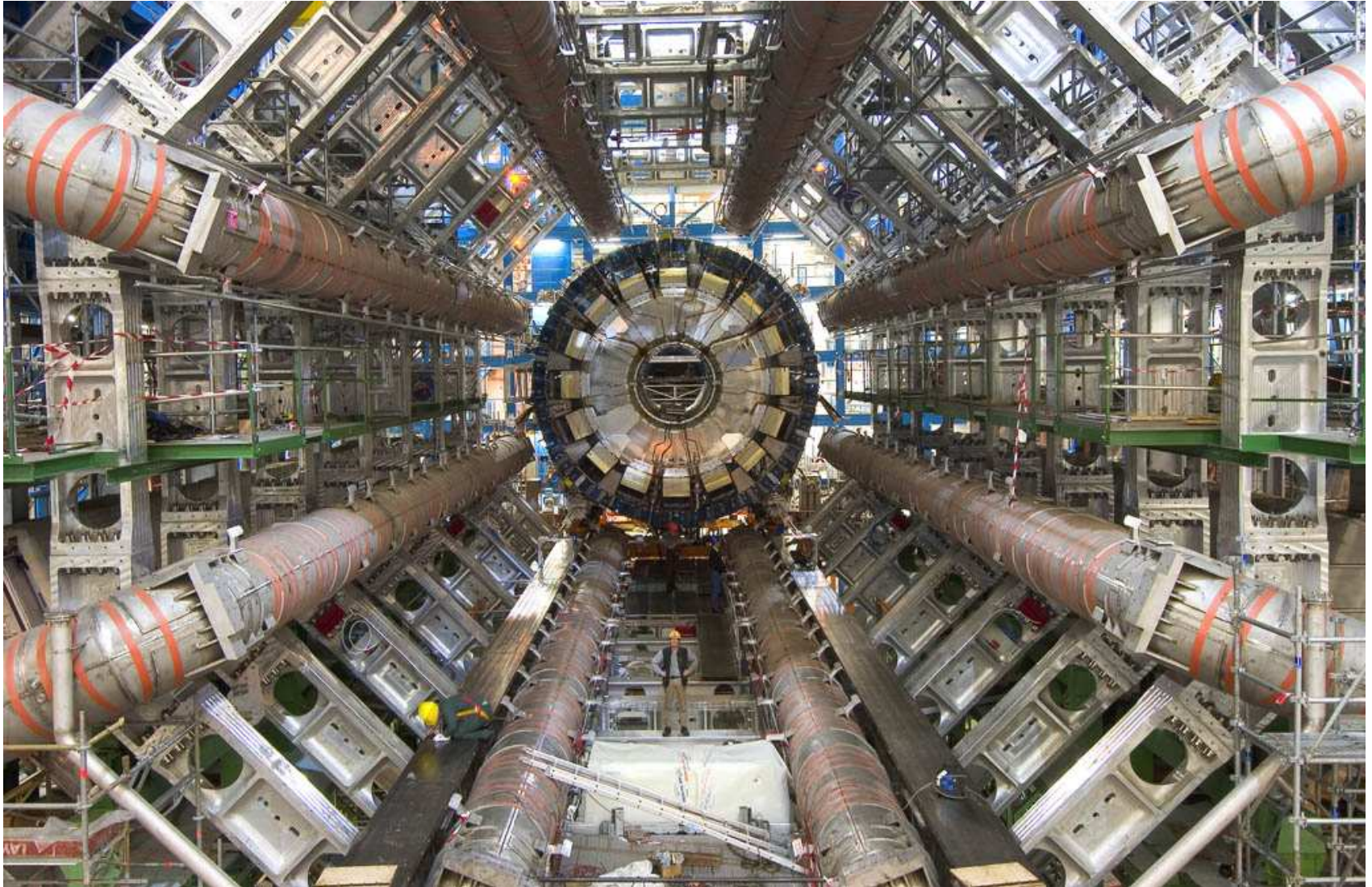
[hep-th/0502058](#) (V. Balasubramanian, P. Berglund, JC, F. Quevedo)

[hep-th/0505076](#) (JC, F. Quevedo, K. Suruliz)

# Talk Structure

- Motivation
- Introduction and Review of Supersymmetry Breaking
- Computing Matter Metrics from Yukawa Couplings
- Results for Particular Models
- Applications: Soft Terms and Neutrino Masses
- Conclusions

# Motivation!



# Motivation

- The LHC!
- This will probe TeV-scale physics in unprecedented detail. If supersymmetry exists, it will (probably) be discovered at the LHC.
- Low-energy supersymmetry represents one of the best possibilities for connecting string theory (or any high-scale theory) to data.
- Understanding supersymmetry breaking and predicting the pattern of superpartners is one of the most important tasks of string phenomenology.

# MSSM Basics

The MSSM is specified by its matter content (that of the supersymmetric Standard Model) plus soft breaking terms. These are

- Soft scalar masses,  $m_i^2 \phi_i^2$
- Gaugino masses,  $M_a \lambda^a \lambda^a$ ,
- Trilinear scalar A-terms,  $A_{\alpha\beta\gamma} \phi^\alpha \phi^\beta \phi^\gamma$
- B-terms,  $BH_1 H_2$ .

It is these soft terms that we want to compute from string compactifications.

# MSSM Basics

- The soft terms break supersymmetry explicitly, but do not reintroduce quadratic divergences into the Lagrangian.
- The LHC will (hopefully) give experimental information about the soft breaking terms and the sparticle spectrum.
- Our task as theorists is to beat the experimentalists to this spectrum!

# SUSY Breaking basics

The soft terms are generated by the mechanism of supersymmetry breaking and how this is transmitted to the observable sector. Examples are

- Gravity mediation - hidden sector supersymmetry breaking
- Gauge mediation - visible sector supersymmetry breaking
- Anomaly mediation - susy breaking through loop effects

These all have characteristic features and scales. Generally,

$$M_{soft} = \frac{M_{susy-breaking}^2}{M_{transmission}}.$$



# SUSY Breaking basics

- In this talk I will focus on gravity mediation.
- This arises naturally in string compactifications, where the hidden sector can be identified with the compactification moduli.
- The formalism for computing soft terms is also well established.

# Gravity Mediation I

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In gravity-mediation, supersymmetry is

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- communicated to the observable sector through non-renormalisable supergravity contact interactions
- which are suppressed by  $M_P$ .

Naively,

$$m_{susy} \sim \frac{F^2}{M_P}$$

This requires  $F \sim 10^{11}$  GeV for TeV-scale soft terms.

# Gravity Mediation II

- The computation of soft terms starts by expanding the supergravity  $K$  and  $W$  in terms of the matter fields  $C^\alpha$ ,

$$W = \hat{W}(\Phi) + \mu(\Phi)H_1H_2 + \frac{1}{6}Y_{\alpha\beta\gamma}(\Phi)C^\alpha C^\beta C^\gamma + \dots,$$

$$K = \hat{K}(\Phi, \bar{\Phi}) + \tilde{K}_{\alpha\bar{\beta}}(\Phi, \bar{\Phi})C^\alpha C^{\bar{\beta}} + [ZH_1H_2 + h.c.] + \dots,$$

$$f_a = f_a(\Phi).$$

- Given this expansion, the computation of the physical soft terms is straightforward.
- The function  $\tilde{K}_{\alpha\bar{\beta}}(\Phi)$  is crucial in computing soft terms, as it determines the normalisation of the matter fields.

# Gravity Mediation III

- Soft scalar masses  $m_{ij}^2$  and trilinears  $A_{\alpha\beta\gamma}$  are given by

$$\begin{aligned} \tilde{m}_{\alpha\bar{\beta}}^2 &= (m_{3/2}^2 + V_0) \tilde{K}_{\alpha\bar{\beta}} \\ &\quad - \bar{F}^{\bar{m}} F^n \left( \partial_{\bar{m}} \partial_n \tilde{K}_{\alpha\bar{\beta}} - (\partial_{\bar{m}} \tilde{K}_{\alpha\bar{\gamma}}) \tilde{K}^{\bar{\gamma}\delta} (\partial_n \tilde{K}_{\delta\bar{\beta}}) \right) \\ A'_{\alpha\beta\gamma} &= e^{\hat{K}/2} F^m \left[ \hat{K}_m Y_{\alpha\beta\gamma} + \partial_m Y_{\alpha\beta\gamma} \right. \\ &\quad \left. - \left( (\partial_m \tilde{K}_{\alpha\bar{\rho}}) \tilde{K}^{\bar{\rho}\delta} Y_{\delta\beta\gamma} + (\alpha \leftrightarrow \beta) + (\alpha \leftrightarrow \gamma) \right) \right]. \end{aligned}$$

- Any physical prediction for the soft terms requires a knowledge of  $\tilde{K}_{\alpha\bar{\beta}}$  for chiral matter fields.
- However,  $\tilde{K}_{\alpha\bar{\beta}}$  is non-holomorphic and thus hard to compute.



# Gravity Mediation IV

To compute soft terms in string compactifications, one must

- choose a particular string compactification.
- compute the moduli potential and determine a susy-breaking minimum.
- evaluate the moduli F-terms at this minimum.
- use the expansions of  $K$  and  $W$  to compute the soft terms.

Note,

- *Any* physical prediction for the soft terms requires a knowledge of  $\tilde{K}_{\alpha\bar{\beta}}$  for chiral matter fields.

# Gravity Mediation V

- In principle this can be done for all types of string compactification - heterotic, type II, M-theory.
- I focus on type IIB compactifications with D3 and D7 branes.
- These are well-studied with regard to the generation of moduli potentials.
- I will mainly discuss techniques to compute the Kähler potential for bifundamental matter fields.

# What is known?

- In Calabi-Yau backgrounds, the Kähler metrics for (non-chiral) D3 and D7 position moduli and D7 Wilson line moduli have been computed by dimensional reduction.

$$\tilde{K}_{D7} \sim \frac{1}{S + \bar{S}} \quad \tilde{K}_{D3} \sim \frac{1}{T + \bar{T}}$$

- Using explicit string scattering computations, matter metrics for bifundamental D7/D7 matter have been computed in IIB toroidal backgrounds.

$$\tilde{K}_{D7_i D7_j} \sim \frac{1}{\sqrt{T_k + \bar{T}_k}}$$

# This talk

- I will describe new techniques for computing the matter metric  $\tilde{K}_{\alpha\bar{\beta}}$  for chiral matter fields on a Calabi-Yau.
- These techniques will apply to chiral bifundamental D7/D7 matter in IIB compactifications.
- I will compute the modular dependence of the matter metrics from the modular dependence of the (physical) Yukawa couplings.

I start by reviewing how Yukawa couplings arise in supergravity.

# Yukawa Couplings in Supergravity

- In supergravity, the Yukawa couplings arise from the Lagrangian terms (Wess & Bagger)

$$\begin{aligned}\mathcal{L}_{kin} + \mathcal{L}_{yukawa} &= \tilde{K}_\alpha \partial_\mu C^\alpha \partial^\mu \bar{C}^{\bar{\alpha}} + e^{\hat{K}/2} \partial_\beta \partial_\gamma W \psi^\beta \psi^\gamma \\ &= \tilde{K}_\alpha \partial_\mu C^\alpha \partial^\mu \bar{C}^{\bar{\alpha}} + e^{\hat{K}/2} Y_{\alpha\beta\gamma} C^\alpha \psi^\beta \psi^\gamma\end{aligned}$$

(This assumes diagonal matter metrics, but we can relax this)

- The matter fields are not canonically normalised.
- The *physical* Yukawa couplings are given by

$$\hat{Y}_{\alpha\beta\gamma} = e^{\hat{K}/2} \frac{Y_{\alpha\beta\gamma}}{(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma)^{\frac{1}{2}}}.$$

# Computing $\tilde{K}_{\alpha\bar{\beta}}$ (I)

$$\hat{Y}_{\alpha\beta\gamma} = e^{\hat{K}/2} \frac{Y_{\alpha\beta\gamma}}{(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma)^{\frac{1}{2}}}.$$

- We know the modular dependence of  $\hat{K}$ :

$$\hat{K} = -2 \ln(\mathcal{V}) - \ln \left( i \int \Omega \wedge \bar{\Omega} \right) - \ln(S + \bar{S})$$

- We compute the modular dependence of  $\tilde{K}_\alpha$  from the modular dependence of  $\hat{Y}_{\alpha\beta\gamma}$ . We work in a power series expansion and determine the leading power  $\lambda$ ,

$$\tilde{K}_\alpha \sim (T + \bar{T})^\lambda k_\alpha(\phi) + (T + \bar{T})^{\lambda-1} k'_\alpha(\phi) + \dots$$

$\lambda$  is the **modular weight** of the field  $T$ .

# Computing $\tilde{K}_{\alpha\bar{\beta}}$ (II)

$$\hat{Y}_{\alpha\beta\gamma} = e^{\hat{K}/2} \frac{Y_{\alpha\beta\gamma}}{(\tilde{K}_{\alpha}\tilde{K}_{\beta}\tilde{K}_{\gamma})^{\frac{1}{2}}}.$$

- We are unable to compute  $Y_{\alpha\beta\gamma}$ . If  $Y_{\alpha\beta\gamma}$  depends on a modulus  $T$ , knowledge of  $\hat{Y}_{\alpha\beta\gamma}$  gives no information about the dependence  $\tilde{K}_{\alpha\bar{\beta}}(T)$ .
- Our results will be restricted to those moduli that do not appear in the superpotential.
- The main example will be the  $T$ -moduli (Kähler moduli) in IIB compactifications. In perturbation theory,

$$\partial_{T_i} Y_{\alpha\beta\gamma} \equiv 0.$$

# Computing $\tilde{K}_{\alpha\bar{\beta}}$ (III)

$$\hat{Y}_{\alpha\beta\gamma} = e^{\hat{K}/2} \frac{Y_{\alpha\beta\gamma}}{(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma)^{\frac{1}{2}}}.$$

- We know  $\hat{K}(T)$ .
- If we can compute  $\hat{Y}_{\alpha\beta\gamma}(T)$ , we can then deduce  $\tilde{K}_\alpha(T)$ .
- Computing  $\hat{Y}_{\alpha\beta\gamma}(T)$  is not as hard as it sounds!
- In IIB compactifications, this can be carried out through wavefunction overlap.

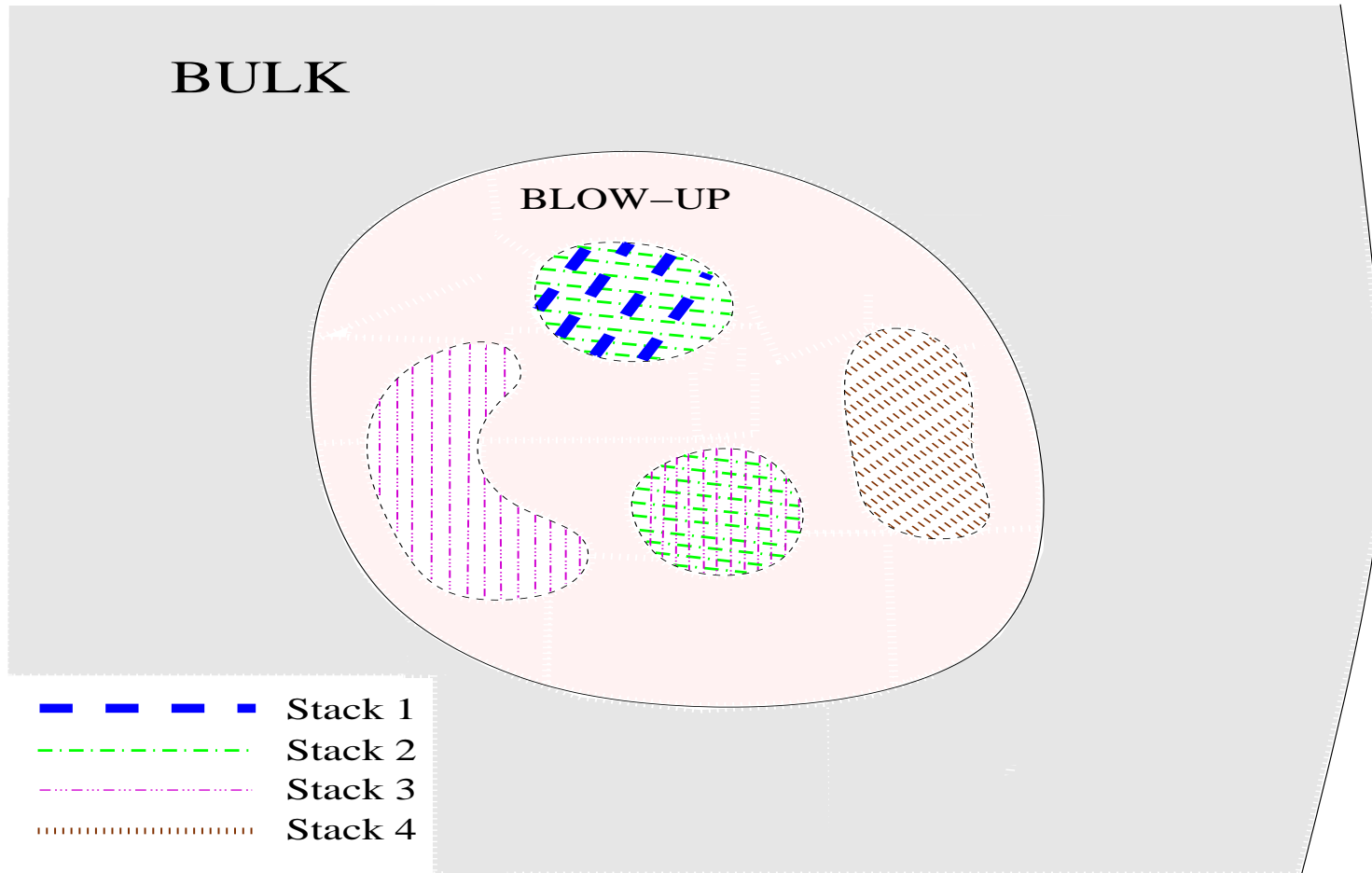
We now describe the computation of  $\hat{Y}_{\alpha\beta\gamma}$  for bifundamental matter on a stack of magnetised D7-branes.



# The brane geometry

- We consider a stack of (magnetised) branes all wrapping an identical cycle.
- If the branes are magnetised, bifundamental fermions can stretch between differently magnetised branes.
- This is a typical geometry in ‘branes at singularities’ models.

# The brane geometry



# Computing $\hat{Y}_{\alpha\beta\gamma}$

- We use a simple computational technique:
- Physical Yukawa couplings are determined by the triple overlap of normalised wavefunctions.
- These wavefunctions can be computed (in principle) by dimensional reduction of the brane action.

# Dimensional reduction

Consider a stack of D7 branes wrapping a 4-cycle  $\Sigma$  in a Calabi-Yau  $X$ . The low-energy limit of the DBI action reduces to super Yang-Mills,

$$S_{SYM} = \int_{\mathbb{M}_4 \times \Sigma} d^n x \sqrt{g} (F_{\mu\nu} F^{\mu\nu} + \bar{\lambda} \Gamma^i (A_i + \partial_i) \lambda)$$

- Magnetic flux on the brane gives bifundamental fermions  $\psi_\alpha$  in the low energy spectrum.
- These fermions come from dimensional reduction of the gaugino  $\lambda_i$ . They are counted topologically by the number of solutions of the Dirac equation

$$\Gamma^i D_i \psi = 0.$$

# Comments

- The full action to be reduced is the DBI action rather than that of Super Yang-Mills.
- In the limit of large cycle volume, magnetic fluxes are diluted in the cycle, and the DBI action reduces to that of super Yang-Mills.
- Our results will hold within this large cycle volume, dilute flux approximation.

# Computing the physical Yukawas $\hat{Y}_{\alpha\beta\gamma}$

Consider the higher-dimensional term

$$\mathcal{L} \sim \int_{\mathbb{M}_4 \times \Sigma} d^n x \sqrt{g} \bar{\lambda} \Gamma^i (A_i + \partial_i) \lambda$$

Dimensional reduction of this gives the kinetic terms

$$\bar{\psi} \partial \psi$$

and the Yukawa couplings

$$\bar{\psi} \phi \psi$$

- The physical Yukawa couplings are set by the combination of the above!

# Kinetic Terms

$$\mathcal{L} \sim \int_{\mathbb{M}_4 \times \Sigma} d^n x \sqrt{g} \bar{\lambda} \Gamma^i (A_i + \partial_i) \lambda$$

- Dimensional reduction gives

$$\lambda = \psi_4 \otimes \psi_6, \quad A_i = \phi_4 \otimes \phi_6.$$

and the reduced kinetic terms are

$$\left( \int_{\Sigma} \psi_6^\dagger \psi_6 \right) \int_{\mathbb{M}_4} \bar{\psi}_4 \Gamma^\mu (A_\mu + \partial_\mu) \psi_4$$

- Canonically normalised kinetic terms require

$$\int_{\Sigma} \psi_6^\dagger \psi_6 = 1.$$

# Yukawa Couplings

$$\mathcal{L} \sim \int_{\mathbb{M}_4 \times \Sigma} d^n x \bar{\lambda} \Gamma^i (A_i + \partial_i) \lambda$$

- Dimensional reduction gives the four-dimensional interaction

$$\left( \int_{\Sigma} \bar{\psi}_6 \Gamma^I \phi_{I,6} \psi_6 \right) \int_{\mathbb{M}_4} d^4 x \phi_4 \bar{\psi}_4 \psi_4.$$

- The physical Yukawa couplings are determined by the (normalised) overlap integral

$$\hat{Y}_{\alpha\beta\gamma} = \left( \int_{\Sigma} \bar{\psi}_6 \Gamma^I \phi_{I,6} \psi_6 \right).$$



# The Result

$$\int_{\Sigma} \psi_6^\dagger \psi_6 = 1, \quad \hat{Y}_{\alpha\beta\gamma} = \left( \int_{\Sigma} \bar{\psi}_6 \Gamma^I \phi_{I,6} \psi_6 \right).$$

- For canonical kinetic terms,

$$\psi_6(y) \sim \frac{1}{\sqrt{\text{Vol}(\Sigma)}}, \quad \Gamma^I \phi_{I,6} \sim \frac{1}{\sqrt{\text{Vol}(\Sigma)}},$$

and so

$$\hat{Y}_{\alpha\beta\gamma} \sim \text{Vol}(\Sigma) \cdot \frac{1}{\sqrt{\text{Vol}(\Sigma)}} \cdot \frac{1}{\sqrt{\text{Vol}(\Sigma)}} \cdot \frac{1}{\sqrt{\text{Vol}(\Sigma)}} \sim \frac{1}{\sqrt{\text{Vol}(\Sigma)}}$$

- This gives the scaling of  $\hat{Y}_{\alpha\beta\gamma}(T)$ .

# Comments

- Q. Under the cycle rescaling, why should  $\psi(y)$  scale simply as

$$\psi(y) \rightarrow \frac{\psi(y)}{\sqrt{\text{Re}(T)}}?$$

A. The  $T$ -moduli do not appear in  $Y_{\alpha\beta\gamma}$  and do not see flavour.

Any more complicated behaviour would alter the form of the triple overlap integral and  $Y_{\alpha\beta\gamma}$  - but this would require altering the complex structure moduli.

- The result for the scaling of  $\hat{Y}_{\alpha\beta\gamma}$  holds in the classical limit of large cycle volume.

# Application: One-modulus KKLT

We can now compute  $\tilde{K}_\alpha(T)$ . In a 1-modulus KKLT model, all chiral matter is supported on D7 branes wrapping the single 4-cycle  $T$ . From above,

$$\hat{Y}_{\alpha\beta\gamma} = e^{\hat{K}/2} \frac{Y_{\alpha\beta\gamma}}{(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma)^{1/2}} \sim \frac{1}{\sqrt{T + \bar{T}}}$$

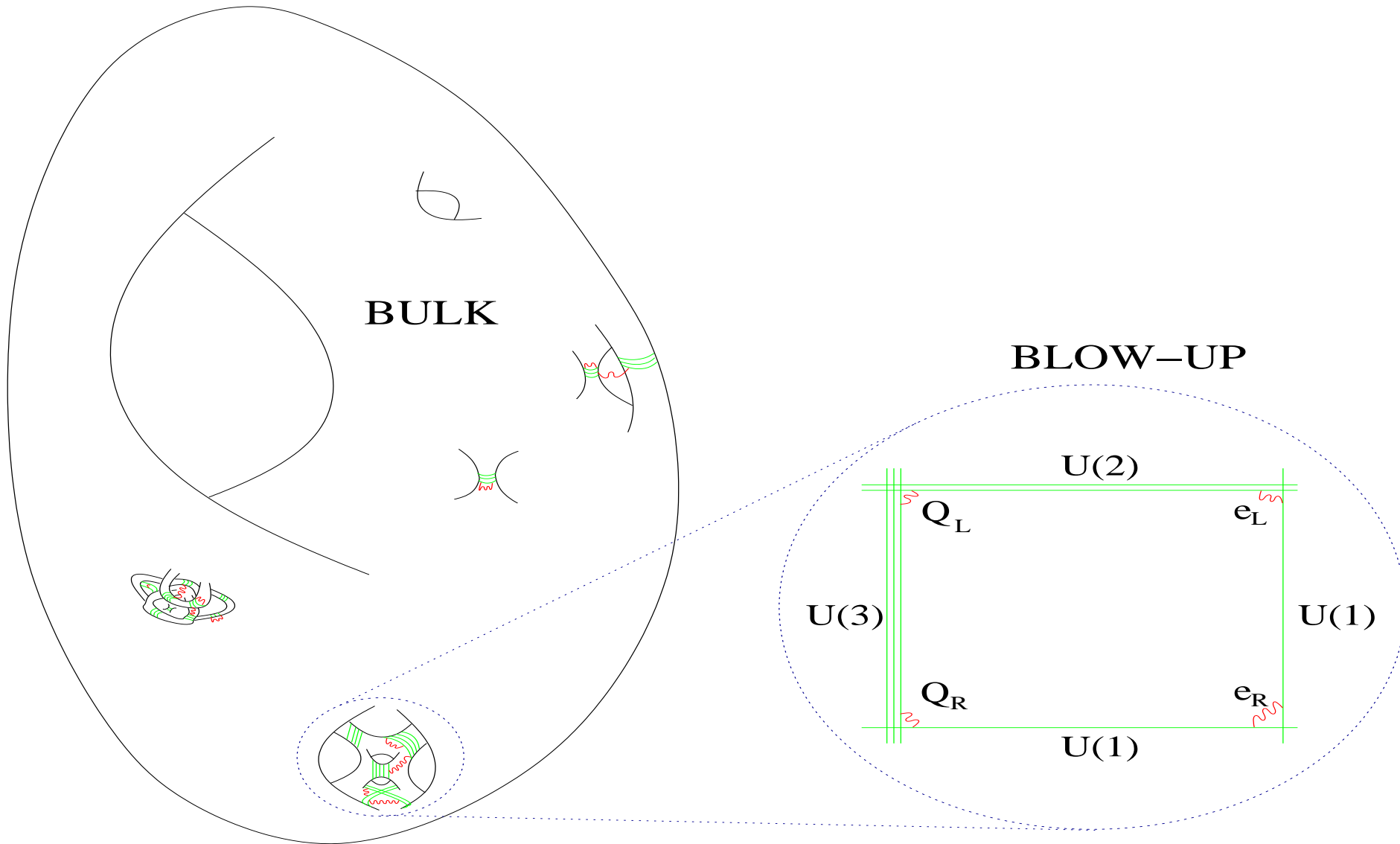
We know the moduli Kähler potential,

$$\hat{K} = -3 \ln(T + \bar{T})$$

and so the matter Kähler potential must scale as

$$\tilde{K}_\alpha \sim \frac{1}{(T + \bar{T})^{2/3}}.$$

# Application: Large-Volume Models



# Application: Large-Volume Models

- These arise in IIB flux compactifications once  $\alpha'$  corrections are included.
- They require  $\geq 2$  Kähler moduli, one 'big' and one 'small'.
- The  $\alpha'$  and non-perturbative corrections compete and determine the structure of the scalar potential.
- The name is because the overall volume is very large,  $\mathcal{V} \sim 10^{15} l_s^6$ , with small cycles  $\tau_s \sim 10 l_s^4$ .

# Application: Large-Volume Models

- For the simplest model  $\mathbb{P}^4_{[1,1,1,6,9]}$ , the Kähler potential is

$$\hat{K} = -2 \ln \mathcal{V} = -2 \ln \left( (T_b + \bar{T}_b)^{3/2} - (T_s + \bar{T}_s)^{3/2} \right).$$

- We can interpret the  $T_s$  cycle as a local, ‘blow-up’ cycle.
- We want  $\tilde{K}_\alpha(T)$  for chiral matter on branes wrapping this cycle.
- The gauge theory supported on a brane wrapping  $T_s$  is determined by local geometry.

# Application: Large-Volume Models

- The physical Yukawa coupling

$$\hat{Y}_{\alpha\beta\gamma} = e^{\hat{K}/2} \frac{Y_{\alpha\beta\gamma}}{(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma)^{\frac{1}{2}}}$$

is local and thus independent of  $\mathcal{V}$ .

- As  $\hat{K} = -2 \ln \mathcal{V}$ , we can deduce *simply from locality* that

$$\tilde{K}_\alpha \sim \frac{1}{\mathcal{V}^{2/3}}.$$

- As this is for a Calabi-Yau background, this is already non-trivial!

# Application: Large-Volume Models

- We can go further. With all branes wrapping the same 4-cycle,

$$\hat{Y}_{\alpha\beta\gamma} = e^{\hat{K}/2} \frac{Y_{\alpha\beta\gamma}}{(\tilde{K}_\alpha \tilde{K}_\beta \tilde{K}_\gamma)^{1/2}} \sim \frac{1}{\sqrt{\tau_s}}.$$

- We can then deduce that

$$\tilde{K}_\alpha \sim \frac{\tau_s^{1/3}}{\mathcal{V}^{2/3}}.$$

- We also have the dependence on  $\tau_s$ !



# Application: Large-Volume Models

- The dependence

$$\tilde{K}_\alpha \sim \frac{1}{\mathcal{V}^{2/3}}$$

follows purely from the requirement that physical Yukawa couplings are local.

- The dependence on  $\tau_s$ ,

$$\tilde{K}_\alpha \sim \frac{\tau_s^{1/3}}{\mathcal{V}^{2/3}}$$

follows from the specific brane configuration (all D7 branes wrapping the same small cycle).

# Application: Large-Volume Models

- If we wrap branes on the large cycle, as for the one-modulus KKLT model we find

$$\tilde{K} \sim \frac{1}{(T_b + \bar{T}_b)^{2/3}}.$$

- All the above results hold for both diagonal and non-diagonal matter metrics.
- The reason is that the  $T$  moduli do not see flavour, and so the  $T$ -dependence is flavour-universal.

# Soft Terms

- We can use the above matter metric to compute the soft terms for the large-volume models.
- We get

$$M_i = \frac{F^s}{2\tau_s} \equiv M,$$

$$m_{\alpha\bar{\beta}} = \frac{M}{\sqrt{3}} \tilde{K}_{\alpha\bar{\beta}},$$

$$A_{\alpha\beta\gamma} = -M \hat{Y}_{\alpha\beta\gamma},$$

$$B = -\frac{4M}{3}.$$

# Soft Terms

- These soft terms are the same as in the dilaton-dominated heterotic scenario.
- They are also flavour-universal.
- This is surprising - there is a naive expectation that gravity mediation will give non-universal soft terms.
- Why? Flavour physics is Planck-scale physics, and in gravity mediation supersymmetry breaking is also Planck-scale physics.
- Naively, we expect the susy-breaking sector to 'see' flavour and thus give non-universal soft terms.

# Soft Terms: Flavour Universality

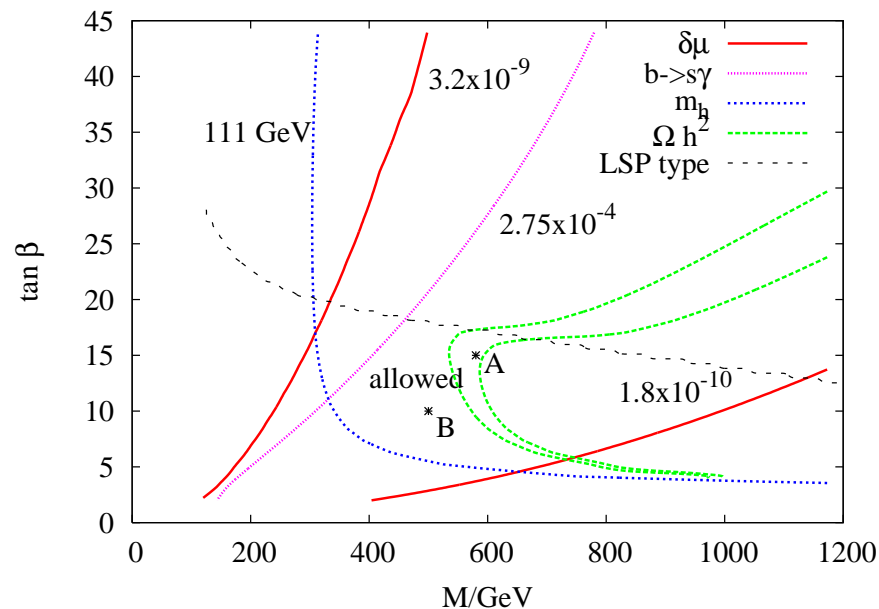
- These expectations are EFT expectations
- In string theory, we have Kähler ( $T$ ) and complex structure ( $U$ ) moduli.
- These are **decoupled** at leading order.

$$\mathcal{K} = -2 \ln(\mathcal{V}(T)) - \ln\left(i \int \Omega \wedge \bar{\Omega}(U)\right) - \ln(S + \bar{S}).$$

- Here,  $U$  sources flavour and  $T$  breaks supersymmetry.
- At leading order, susy-breaking and flavour decouple.
- The origin of universality is the decoupling of Kähler and complex structure moduli.

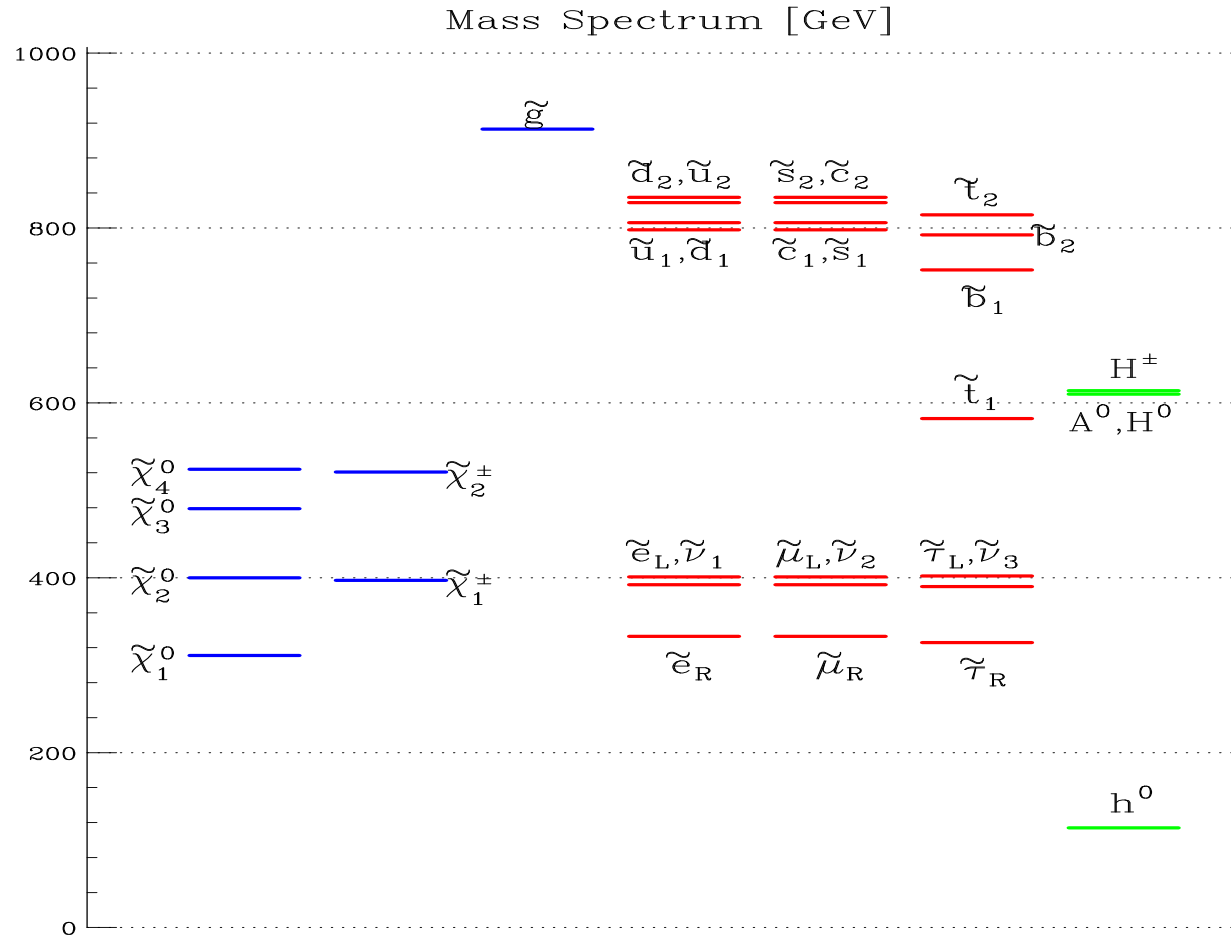
# Soft Terms: Phenomenology

- We run the soft terms to low energy using SoftSUSY and study the spectrum.
- The constraints are given by



# Soft Terms: Phenomenology

Spectrum:



# Neutrino Masses

- The theoretical origin of neutrino masses is a mystery. Experimentally

$$0.05\text{eV} \lesssim m_\nu^H \lesssim 0.3\text{eV}.$$

- This corresponds to a Majorana mass scale

$$M_{\nu_R} \sim 3 \times 10^{14}\text{GeV}.$$

- Equivalently, this is the suppression scale  $\Lambda$  of the dimension five Standard Model operator

$$\mathcal{O} = \frac{1}{\Lambda} H H L L.$$



# Neutrino Masses

- In the supergravity MSSM, consider the superpotential operator

$$\frac{\lambda}{M_P} H_2 H_2 L L \in W,$$

where  $\lambda$  is dimensionless.

- This corresponds to the *physical* coupling

$$e^{\hat{K}/2} \frac{\lambda}{M_P} \frac{\langle H_2 H_2 \rangle L L}{(\tilde{K}_{H_2} \tilde{K}_{H_2} \tilde{K}_L \tilde{K}_L)^{\frac{1}{2}}}.$$

- Once the Higgs receives a vev, this generates neutrino masses.

# Neutrino Masses

- Using the large-volume result  $\tilde{K}_\alpha \sim \frac{\tau_s^{1/3}}{\mathcal{V}^{2/3}}$ , this becomes

$$\frac{\lambda \mathcal{V}^{1/3}}{\tau_s^{2/3} M_P} \langle H_2 H_2 \rangle LL.$$

- Use  $\mathcal{V} \sim 10^{15}$  (to get  $m_{3/2} \sim 1\text{TeV}$ ) and  $\tau_s \sim 10$ :

$$\frac{\lambda}{10^{14} \text{GeV}} \langle H_2 H_2 \rangle LL$$

- With  $\langle H_2 \rangle = \frac{v}{\sqrt{2}} = 174\text{GeV}$ , this gives

$$m_\nu = \lambda(0.3 \text{ eV}).$$

# Large Volumes are Powerful

- In large-volume models, an exponentially large volume appears naturally ( $\mathcal{V} \sim e^{\frac{c}{g_s}}$ ).
- A volume  $\mathcal{V} \sim 10^{15} l_s^6$ , required for TeV-scale supersymmetry, automatically gives the correct scale for neutrino masses.
- This same volume also gives an axion decay constant in the allowed window (JC, [hep-th/0602233](#)),

$$10^9 \text{ GeV} \lesssim f_a \lesssim 10^{12} \text{ GeV}.$$

- This yoking of three distinct scales is very attractive.
- The origin of all three hierarchies is the exponentially large volume.

# Conclusions

- Kähler metrics for bifundamental matter enter crucially in the computation of MSSM soft terms.
- I have described techniques to compute these in IIB Calabi-Yau string compactifications.
- We computed the modular weights both for one-modulus KKLT models and for the multi-modulus large-volume models.
- The soft terms are flavour-universal, which comes from the decoupling of Kähler and complex structure moduli.
- For the large-volume models, TeV-scale supersymmetry naturally gives the correct scale for neutrino masses.

