Phenomenological Aspects of LARGE Volume Models

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Shehu Abdussalam, Ben Allanach, Vijay Balasubramanian, Per Berglund, Cliff Burgess, Michele Cicoli, Daniel Cremades, Ling-Han Hung, Chun-Hay (Steve) Kom, Fernando Quevedo, Kerim Suruliz

This talk is particularly based on the papers

hep-th/0602233, J. P. Conlon

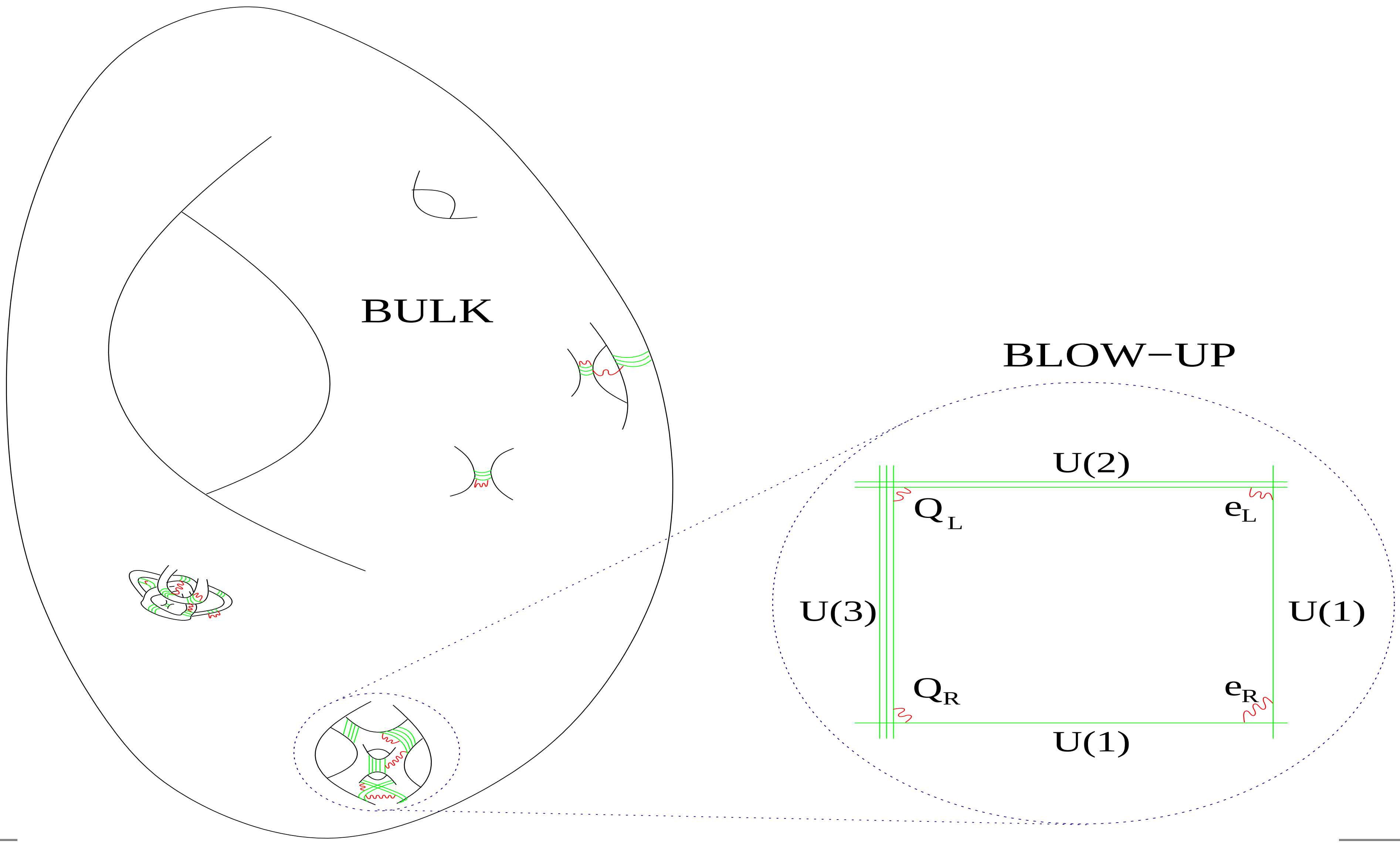
hep-ph/0611144, J. P. Conlon, D. Cremades

0805.xxxx (hep-th), C. P. Burgess, J. P. Conlon, L-H. Hung, C. Kom, A. Maharana, F. Quevedo

Talk Structure

1. TeV Supersymmetry
2. The QCD Axion and the Strong CP Problem
3. Continuous Global Flavour Symmetries
4. Neutrino Masses
5. Hyper-Weak Gauge Groups

Recap of LARGE Volume Models



Recap of LARGE Volume Models

- The stabilised volume is naturally exponentially large.
- The Calabi-Yau has a ‘Swiss cheese’ structure.
- There is a large bulk cycle and a small blow-up cycle.
- The LARGE volume lowers the gravitino mass through

$$m_{3/2} = \frac{M_P W_0}{\mathcal{V}}.$$

- The LARGE volume is a powerful tool to generate various hierarchies - this talk.

TeV Supersymmetry

Supersymmetry is broken in the vacuum.

The supersymmetry breaking scale is set by the gravitino mass

$$M_{3/2} = \frac{W_0 M_P}{\mathcal{V}}$$

The stabilised volume is

$$\mathcal{V} \sim W_0 e^{\frac{c}{g_s}}, \qquad (c \text{ constant})$$

and is essentially arbitrary.

We like TeV supersymmetry and so set $\mathcal{V} \sim 10^{15}$ and explore the consequences!

Axions

Axions are a well-motivated solution to the strong CP problem of the Standard Model.

The Lagrangian for QCD contains a term

$$\mathcal{L}_{QCD} = \frac{1}{4\pi g^2} \int d^4x F^a_{\mu\nu} F^{a,\mu\nu} + \frac{\theta}{8\pi^2} \int F^a \wedge F^a.$$

The θ term

$$\mathcal{L}_\theta = \frac{\theta}{8\pi^2} \int F^a \wedge F^a.$$

gives rise to the strong CP problem.

This term vanishes in perturbation theory and is only non-vanishing for non-perturbative instanton computations.

Axions

θ is an angle and in principle takes any value $\theta \in (-\pi, \pi)$.

It can be shown that non-vanishing θ generates an electric dipole moment for the neutron.

The measured absence of a neutron electric dipole moment implies that

$$|\theta_{QCD}| \lesssim 10^{-10}.$$

No symmetry of the Standard Model requires $|\theta|$ to be so small.

The problem of why $|\theta_{QCD}| \ll \pi$ is the *strong CP* problem of the Standard Model.

Axions

The best solution to the strong CP problem is due to Peccei and Quinn.

The angle θ is promoted to a dynamical field with Lagrangian

$$\mathcal{L} = \frac{1}{2}\partial_\mu\theta\partial^\mu\theta + \int \frac{\theta}{8\pi^2 f_a} F^a \wedge F^a.$$

f_a is called the *axion decay constant* and measures the strength of the non-renormalisable axion-matter coupling.

Axions

Non-perturbative QCD effects generate a potential for θ,

$$V_\theta \sim \Lambda_{QCD}^4 \left(1 - \cos\left(\frac{\theta}{f_a}\right)\right) \quad = \quad \frac{\Lambda_{QCD}^4}{2f_a^2}(\theta^2 + \ldots)$$

The axion field θ gets a mass

$$m_\theta = \frac{\Lambda_{QCD}^2}{f_a} \sim 0.01\text{eV}\left(\frac{10^9\text{GeV}}{f_a}\right)$$

Axion phenomenology depends entirely on the value of f_a.

Axions

Constraints from supernova cooling and direct searches imply $f_a \gtrsim 10^9$GeV (axions cannot couple too strongly to matter).

Avoiding the overproduction of axion dark matter during a hot big bang prefers $f_a \lesssim 10^{12}$GeV (axions cannot couple too weakly to matter).

There exists an axion 'allowed window',

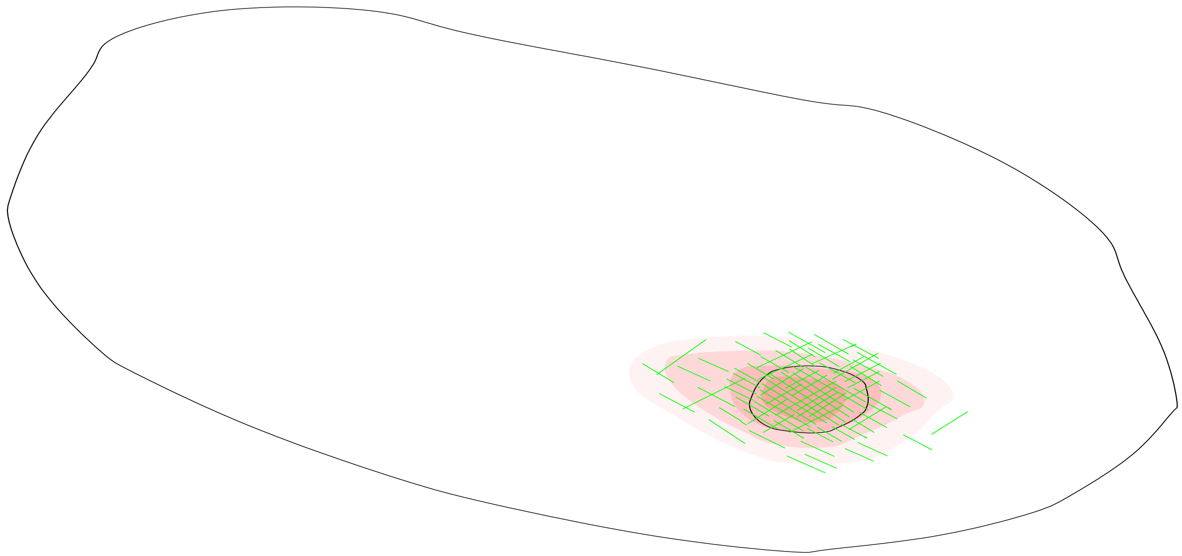
$$10^9\text{GeV} \lesssim f_a \lesssim 10^{12}\text{GeV}.$$

Axions

- In string theory, the axionic coupling comes from the Chern-Simons interaction

$$\int_{\mathbb{M}_4 \times \Sigma_4} F \wedge F \wedge C_4 \to \int_{\mathbb{M}_4} \frac{c}{f_a} F \wedge F.$$

- f_a measures the axion-matter coupling.



Axions

- The axion coupling is a local coupling and does not see the overall volume.
- The coupling is set by the string scale:

$$f_a \sim m_s \sim \frac{M_P}{\sqrt{\mathcal{V}}}.$$

- This generates the axion scale,

$$f_a \sim \frac{M_P}{\sqrt{\mathcal{V}}} \sim 10^{11}\text{GeV}.$$

Axions

TeV-scale supersymmetry required

$$M_{susy} \sim m_{3/2} = \frac{M_P}{\mathcal{V}} = 1\text{TeV}.$$

This gives $\mathcal{V} \sim 10^{15} l_s^6$.

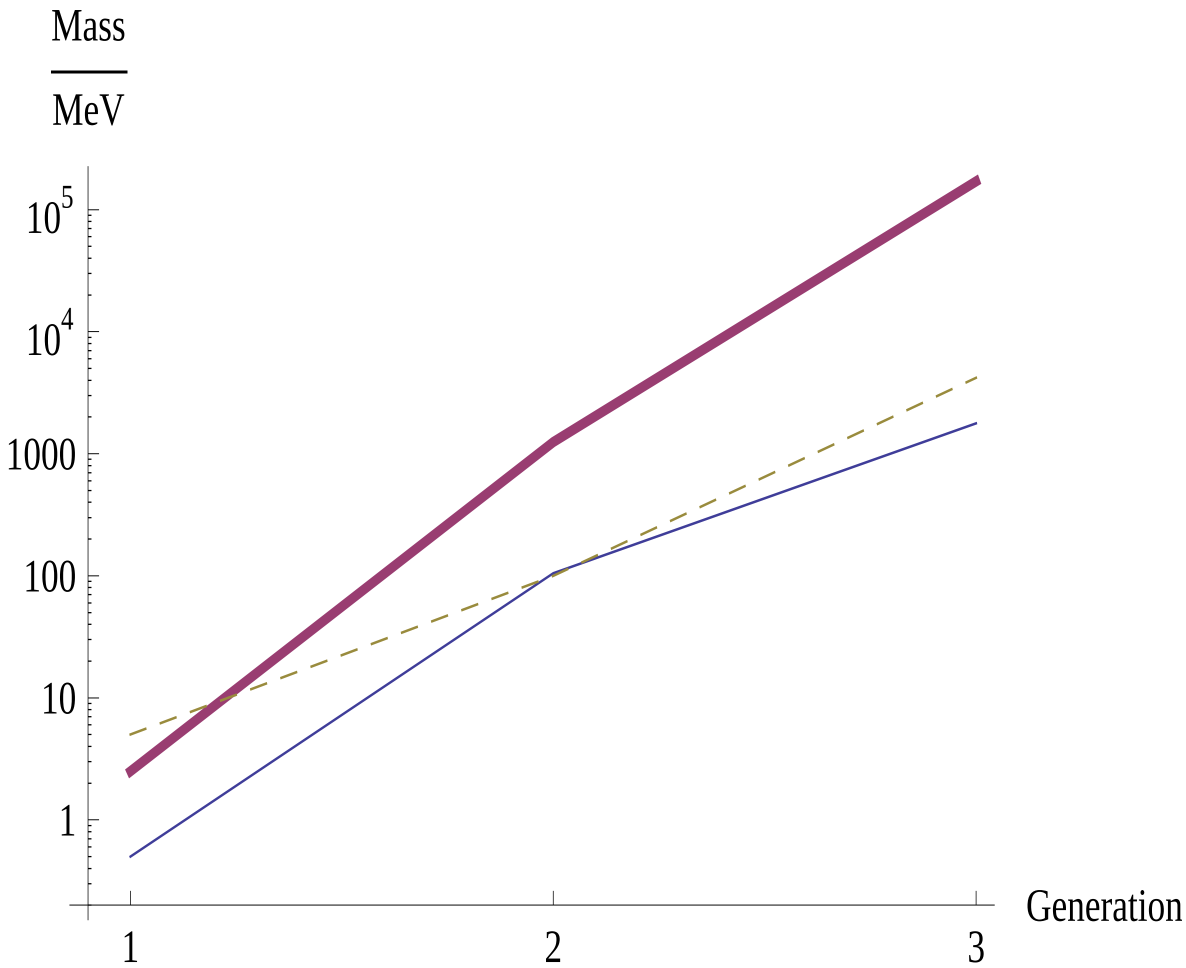
The axion scale is therefore

$$f_a \sim \sqrt{M_{susy} M_P} \sim 10^{11}\text{GeV}$$

The existence of an axion in the allowed window correlates to the existence of supersymmetry at the TeV scale.

Continuous Flavour Symmetries

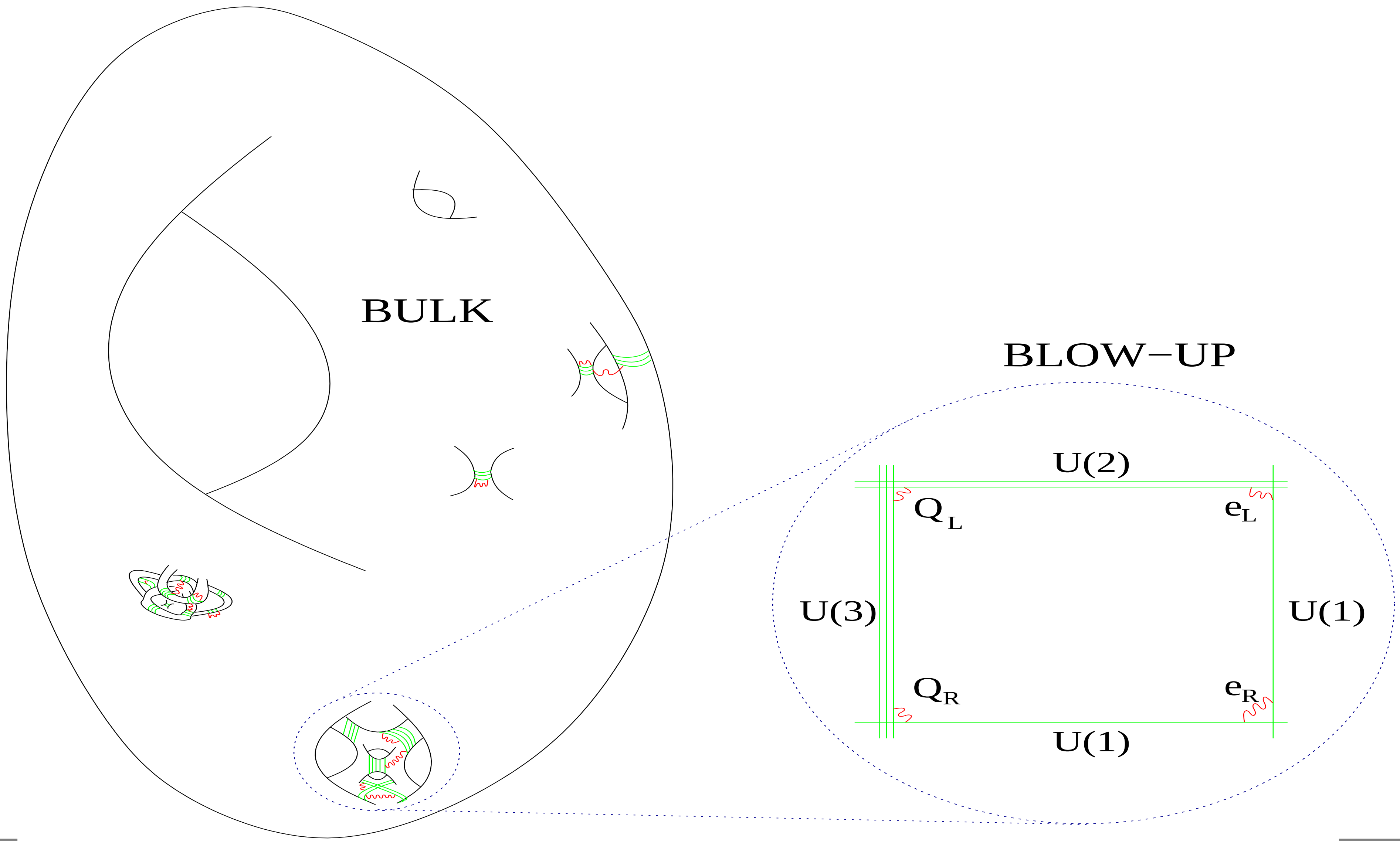
Flavour symmetries are very attractive for explaining the fermion mass hierarchies.



Continuous Flavour Symmetries

- In LARGE volume models, the Standard Model is necessarily a local construction.
- The couplings of the Standard Model are determined by the local geometry and are insensitive to the bulk.
- In the limit $\mathcal{V} \to \infty$, the bulk decouples and all couplings and interactions of the Standard Model are set by the local geometry and metric.
- Global Calabi-Yau metrics are hard - local metrics are known!

Continuous Flavour Symmetries



Continuous Flavour Symmetries

It is a theorem that compact Calabi-Yaus have no isometries.

Local Calabi-Yau metrics often have isometries. Examples:

1. Flat space $\mathbb{C}^3$ has an $SO(6)$ isometry.
2. The (resolved) orbifold singularity $\mathbb{C}^3/\mathbb{Z}_3 = \mathcal{O}_{\mathbb{P}^2}(-3)$ has an $SU(3)/\mathbb{Z}_3$ isometry.
3. The conifold geometry $\sum z_i^2 = 0$ has an $SU(2) \times SU(2) \times U(1)$ isometry.

Local metric isometries are flavour symmetries of local brane constructions.

Continuous Flavour Symmetries

- There are two scales in the geometry - the size of the local metric (set by τ_s) and the size of the global metric (set by τ_b).
- The rescaling $\tau_s \to \lambda^4 \tau_s, \tau_b \to \lambda^4 \tau_b$ rescales the overall metric and is simply a scale factor.
- The presence and goodness of the isometry is set by the ratio $\frac{\tau_s}{\tau_b}$.
- This determines the extent to which the local non-compact metric is a good approximation in the compact case.

Continuous Flavour Symmetries

- In the limit $\mathcal{V} \to \infty$ the flavour symmetry becomes exact and the space becomes non-compact.

  New massless states exist as the bulk KK modes become massless.

- In the limit of $\mathcal{V} \gg 1$ but finite, the flavour symmetry is approximate, being softly broken.

  The bulk KK modes are massive but are hierarchially lighter than the local KK modes:

$$M_{KK,local} = \frac{M_S}{R_s}, \qquad M_{KK,bulk} = \frac{M_S}{R_b}.$$

  The bulk KK modes represent the pseudo-Goldstone bosons of the approximate flavour symmetry.

Continuous Flavour Symmetries

Phenomenological discussions of flavour symmetries start with a symmetry group

$$G_{SM} \times G_F = G_{SU(3)\times SU(2)\times U(1)} \times G_F.$$

Flavons Φ are charged under G_F and not under G_{SM}.
SM matter C_i is charged under both G_F and G_{SM}.

$$W = (\Phi_\alpha \Phi_\beta \Phi_\gamma \ldots) C_i C_j C_k.$$

Flavon vevs break G_F and generate Yukawa textures.
The order parameter for G_F breaking is $< \Phi >$.

What are the flavons in our case???

Continuous Flavour Symmetries

A puzzle:

- In 4d effective theory, flavour symmetry breaking is parametrised by the ratio $\frac{R_s}{R_b} = \frac{\tau_s^{1/4}}{\tau_b^{1/4}}$.

  This sets the relative size of the bulk and local cycles.
- However $\frac{\tau_s^{1/4}}{\tau_b^{1/4}}$ is a real singlet while the flavour symmetry group is non-Abelian.
- $\frac{\tau_s^{1/4}}{\tau_b^{1/4}}$ can only be in a trivial representation of G_F.
- So there are no flavons in the 4d effective field theory!

Continuous Flavour Symmetries

The resolution:

- There are indeed no flavons in the 4d effective field theory!
- The aproximate isometry comes from the full Calabi-Yau metric.
- The flavon modes that are charged under G_F are the higher-dimensional (Kaluza-Klein) modes.

  Yau's theorem implies that the vevs of KK modes are set by the moduli vevs.

  From a 4d perspective, it is the vevs of KK modes that break the flavour symmetry.

Continuous Flavour Symmetries

- This mechanism describes a way to realise continuous global non-Abelian flavour symmetries in string theory.
- (Approximate) flavour symmetries arise from (approximate) isometries of the compact Calabi-Yau.
- The symmetries are exact at infinite volume and are broken for any finite value of the volume.
- The breaking parameter is (R_s/R_b) : LARGE volume implies a small breaking parameter for the flavour symmetry.
- Work in progress: what is the precise breaking parameter within the Yukawa couplings?

Neutrino Masses

- Neutrino masses exist:

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

- In the seesaw mechanism, this corresponds to a Majorana mass scale for right-handed neutrinos

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

- Equivalently, this is the suppression scale Λ of the dimension five MSSM operator

$$\begin{aligned} \mathcal{O}_{m_\nu} &= \frac{1}{\Lambda} H_2 H_2 L L \\ \Rightarrow m_\nu &= 0.1\text{eV} \left(\sin^2\beta \times \frac{3 \times 10^{14}\text{GeV}}{\Lambda} \right). \end{aligned}$$

Neutrino Masses

Neutrino masses imply a scale $\Lambda \sim (\text{a few}) \times 10^{14}$GeV which is

- not the Planck scale 10^{18}GeV
- not the GUT scale 10^{16}GeV
- not the intermediate scale 10^{11}GeV
- not the TeV scale 10^{3}GeV

Can the intermediate-scale string give a quantitative understanding of this scale?

Neutrino Masses

The low-energy theory is

$$W = Y_{\alpha\beta\gamma} C^\alpha C^\beta C^\gamma + \frac{Z_{\alpha\beta\gamma\delta}}{M_P} C^\alpha C^\beta C^\gamma C^\delta + \frac{\lambda}{M_P} H_2 H_2 L L,$$

$$K = \hat{K}(\Phi, \bar{\Phi}) + \tilde{K}_{\alpha\bar{\beta}} \Phi^\alpha \bar{\Phi}^{\bar{\beta}} + \dots$$

We focus particularly on the term

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

This term generates neutrino masses when the Higgs gets a vev.

Neutrino Masses

The physical normalised operator is

$$e^{\hat{K}/2}\frac{\lambda}{M_P}\frac{H_2H_2LL}{\left(\tilde{K}_{H\bar{H}}\tilde{K}_{H\bar{H}}\tilde{K}_{L\bar{L}}\tilde{K}_{L\bar{L}}\right)^{\frac{1}{2}}}$$

We know

$$\hat{K} = -2\ln(\mathcal{V})$$

To compute the physical suppression scale, we need the volume scaling of $\tilde{K}_{H\bar{H}}$ and $\tilde{K}_{L\bar{L}}$.

Neutrino Masses

We can compute this using the Yukawa couplings.

The physical Yukawa couplings are

$$\hat{Y}_{\alpha\beta\gamma} = e^{\hat{K}/2} \frac{Y_{\alpha\beta\gamma}}{\sqrt{K_{\alpha\bar{\alpha}} K_{\beta\bar{\beta}} K_{\gamma\bar{\gamma}}}}$$

Geometry implies these couplings are local.

We know

- $Y_{\alpha\beta\gamma}$ does not depend on $\mathcal{V}$.
- The overall term $e^{\hat{K}/2} \sim \frac{1}{\mathcal{V}}$.

Neutrino Masses

The local interactions care only about the local geometry and decouple from the bulk volume.

Physical locality then implies the physical Yukawa couplings $\hat{Y}_{\alpha\beta\gamma}$ do not depend on the bulk volume.

For local fields C^α the relation

$$\hat{Y}_{\alpha\beta\gamma} = e^{\hat{K}/2} \frac{Y_{\alpha\beta\gamma}}{\sqrt{K_{\alpha\bar{\alpha}} K_{\beta\bar{\beta}} K_{\gamma\bar{\gamma}}}}$$

implies

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

Neutrino Masses

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

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

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

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

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

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

Neutrino Masses

Integrating out string / KK states generates a dimension-five operator suppressed by

$$\text{(string)} \qquad \mathcal{V}^{-1/12} \times \frac{\mathcal{V}^{\frac{1}{2}}}{M_P} \times \mathcal{V}^{-1/12} \sim \frac{\mathcal{V}^{1/3}}{M_P}$$

$$\text{(KK)} \qquad \mathcal{V}^{-1/6} \times \frac{\mathcal{V}^{2/3}}{M_P} \times \mathcal{V}^{-1/6} \sim \frac{\mathcal{V}^{1/3}}{M_P}$$

- Integrating out heavy states of mass M does not produce operators suppressed by M^{-1}.
- The dimension-five suppression scale is independent of the masses of the heavy states integrated out.

Hyper-Weak Gauge Groups

- In LARGE volume models, the Standard Model is a local construction and branes only wrap small cycles.
- There are also bulk cycles associated to the overall volume. These have cycle size

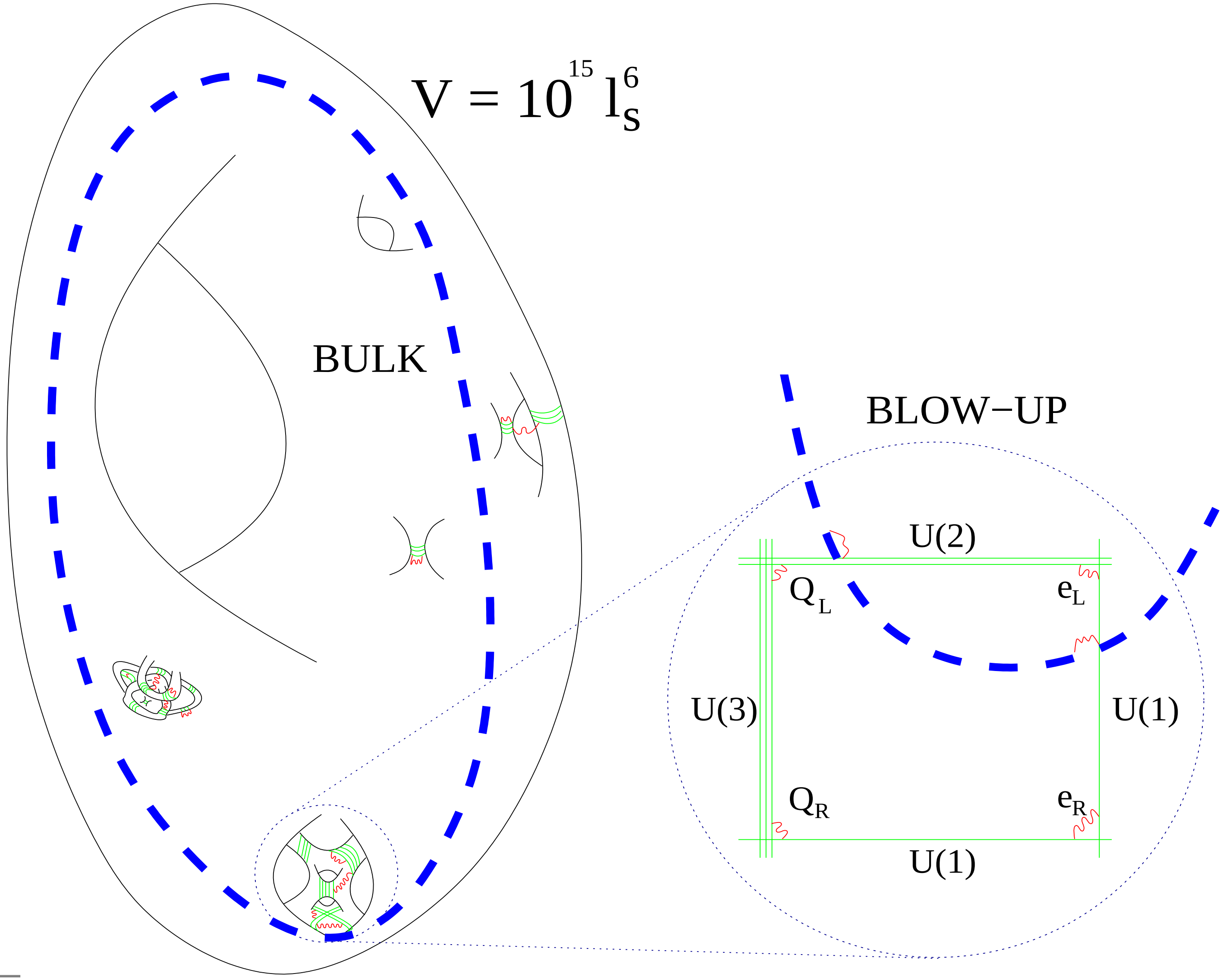
$$\tau_b \sim \mathcal{V}^{2/3} \sim 10^{10}.$$

- There is no reason not to have D7 branes wrapping these cycles!
- The gauge coupling for such branes is

$$\frac{4\pi}{g^2} = \tau_b$$

with $g \sim 10^{-4}$.

Hyper-Weak Gauge Groups



Hyper-Weak Gauge Groups

In LARGE volume models it is a generic expectation that there will exist additional gauge groups with very weak coupling

$$\alpha^{-1} \sim 10^9.$$

Two phenomenological questions to ask:

1. How heavy is the hyper-weak Z' gauge boson?
2. How does Standard Model matter couple to the hyper-weak force?

Hyper-Weak Gauge Groups

- Standard Model can couple directly to bulk D7.
- If weak-scale vevs of H_1, H_2 break the bulk D7 gauge group, then

$$M_{Z'} \sim gv \sim 10^{-4} \times 100\text{GeV} \sim 10\text{MeV}.$$

- Gauge group may couple directly to electrons, but with

$$g_e \sim 10^{-4}, \qquad g_e^2 \sim 10^{-8}.$$

- In this case we have a new (relatively) light very weakly coupled gauge boson.

Hyper-Weak Gauge Groups

- There may also be kinetic mixing between the new gauge boson and the photon (cf Z/γ mixing).
- If the new gauge boson is light, the mixing allows the new boson to couple to electromagnetic currents:

$$\mathcal{L}_{int} = \frac{(\bar{\psi}\gamma^{\mu}\psi)A_{\mu}}{\sqrt{1-\lambda^2}} - \frac{(\bar{\psi}\gamma^{\mu}\psi)\lambda Z'_{\mu}}{\sqrt{1-\lambda^2}}$$

  where λ is the mixing parameter.
- This can give milli-charged fermions under the new gauge boson Z'.

Hyper-Weak Gauge Groups

- In string theory the presence of gauge groups with very small intrinsic coupling can only occur if $M_S \ll M_{planck}$.
- (This is related to the statement that 'gravity is the weakest force').
- Any observation of such gauge groups would be a distinctive signature of LARGE volume physics.

Conclusions

- Moduli stabilisation is an important problem in string theory.
- LARGE volume models are an attractive method of moduli stabilisastion.
- These models have interesting consequences for
  1. Low-energy supersymmetry
  2. QCD axions in the allowed window
  3. The scale of neutrino masses
  4. Realising Standard Model flavour symmetries
  5. Possible new hyper-weak gauge groups