From Strong CP to the LHC - via the Landscape

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String Vacua and the Landscape
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- String theory has many vacua.
- This is a problem.
- How can we choose between vacua?
- How can we be predictive?

- We would like a 'magic bullet' to pick out the true vacuum.

  This does not (currently) exist.

- Instead, use what we know about low energy physics.

  In this talk I will use the strong CP problem.

- I will make one main assumption:

A (Peccei-Quinn) string theory axion solves the strong CP problem.

and use it as a gigantic lever to tell us about

- the landscape
- moduli stabilisation
- supersymmetry breaking
- the sparticle spectrum

I. Introduction

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1. The Strong CP problem
2. Axions in String Theory
3. Axions and Moduli Stabilisation
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III. From the Landscape to the LHC

5. Review of soft masses
6. Moduli stabilisation and soft masses

7. Conclusions (again)

1. The Strong CP Problem

- The strong interactions ought to violate CP, but they do not.

$$\frac{1}{4g^2}\int F^a_{\mu\nu}F^{a\mu\nu} + \underbrace{\frac{\theta}{16\pi^2}\int F^a_{\mu\nu}\tilde{F}^{a\mu\nu}}^{\text{CP violating}}$$

$$\left(\tilde{F}^{\mu\nu} = \epsilon^{\mu\nu\rho\sigma}F_{\rho\sigma}\right)$$

- The CP violating term $\frac{\theta}{16\pi^2}\int F^a_{\mu\nu}\tilde{F}^{a\mu\nu}$ should be present, but experimentally

$$|\theta| < 10^{-10}$$

This is the strong CP problem.

There exist 3 proposed solutions

① $m_u = 0$ and Θ can be rotated away.

- disfavoured by lattice data

② CP is an exact symmetry at high scales which is spontaneously broken (Nelson-Barr)

③ A Peccei-Quinn axion sets $\Theta = 0$ dynamically.

This talk is about ③

The Peccei-Quinn approach:

- In $\frac{\theta}{16\pi^2}\int F^a_{\mu\nu}\tilde{F}^{a\,\mu\nu}$, promote θ to a dynamical field $\theta(x)$ with Lagrangian

$$\mathcal{L} = \mathcal{L}_{SM} + \tfrac{1}{2} f_a^2 \, \partial_\mu\theta \, \partial^\mu\theta + \frac{\theta}{16\pi^2} F_{\mu\nu}\tilde{F}^{\mu\nu}$$

f_a has dimensions of mass and is the axion decay constant.

Canonically normalise:

$$a \equiv \theta f_a$$

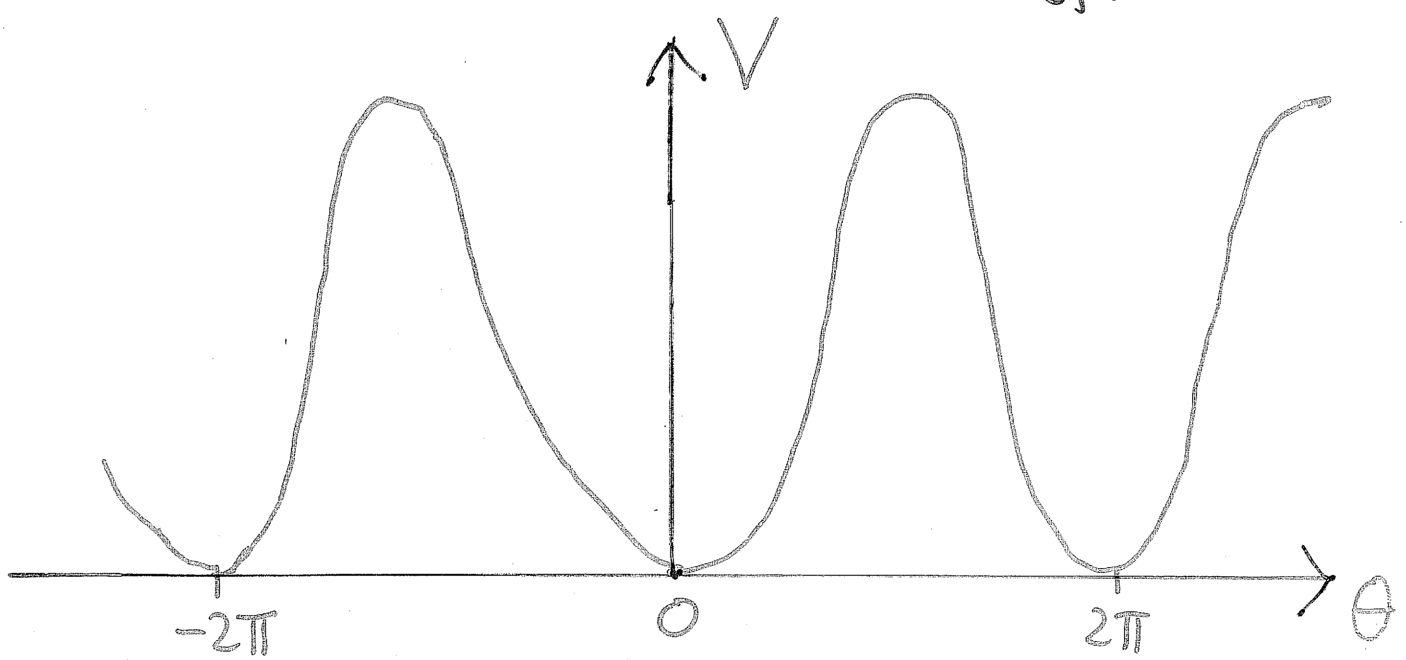
$$\Rightarrow \mathcal{L} = \mathcal{L}_{SM} + \tfrac{1}{2}\partial_\mu a \, \partial^\mu a + \frac{1}{16\pi^2}\frac{a}{f_a} F^a_{\mu\nu}\tilde{F}^{a\,\mu\nu}$$

a is an axion.

The θ angle is now dynamically determined by the potential for a.

QCD instanton effects generate a potential

$$V(a) \sim \Lambda_{QCD}^4 \left(1 - \cos\left(\frac{a}{f_a}\right)\right)$$



This dynamically minimises $\theta = \frac{a}{f_a}$ at 0 and solves the strong CP problem.

The axion Lagrangian

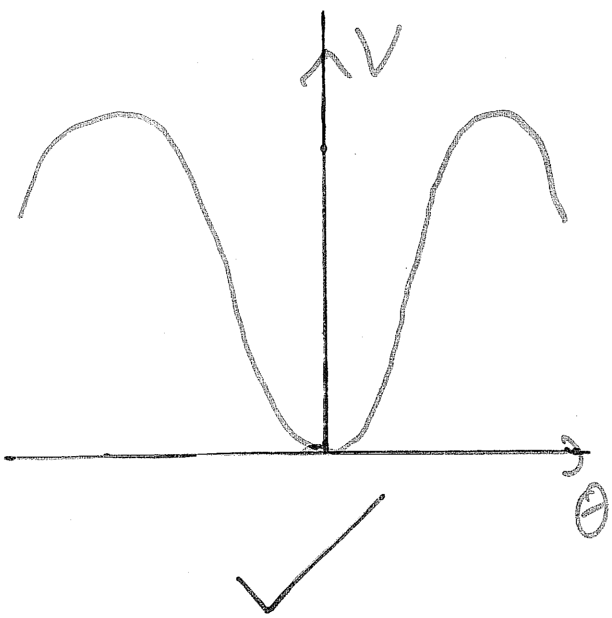
$$\mathcal{L} = \mathcal{L}_{SM} + \frac{1}{2}\partial_\mu a\, \partial^\mu a + \frac{1}{16\pi^2}\frac{a}{f_a} F^a_{\mu\nu}\tilde{F}^{a\mu\nu}$$

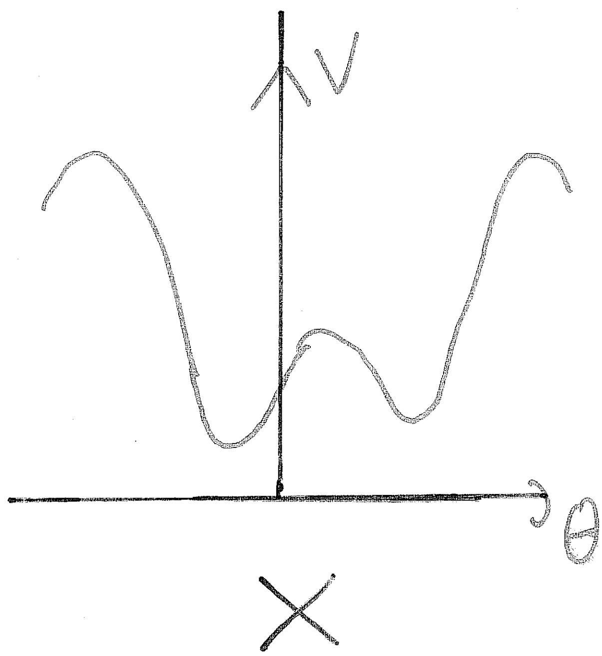
has a $U(1)$ symmetry

$$a \to a + \epsilon$$

broken by QCD instantons. Such a symmetry is a Peccei-Quinn symmetry.

The Peccei-Quinn solution to the strong CP problem requires that QCD instantons dominate the breaking of this symmetry: otherwise, $V(\theta)$ is not minimised at $\theta = 0$.





In standard cosmology, the decay constant f_a is tightly constrained:

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

The lower bound comes from astrophysics: axions do not contribute to supernova cooling.

The upper bound is cosmological: axions do not overclose the universe. This may be relaxed with non-standard cosmologies.

Summary of the Strong CP Problem

- $\mathcal{L} \sim \mathcal{L}_{SM} + \frac{1}{4g^2}\int F^a_{\mu\nu} F^{a\mu\nu} + \frac{\theta}{16\pi^2}\int F^a_{\mu\nu} \tilde{F}^{a\mu\nu}$

  with $|\theta| < 10^{-10}$.

- Solve by

$$\mathcal{L} \sim \mathcal{L}_{SM} + \frac{1}{2}\partial_\mu a \partial^\mu a + \frac{1}{16\pi^2}\int \frac{a}{f_a} \tilde{F}^a_{\mu\nu} F^{a\mu\nu}$$

- The dominant contribution to the potential for a must come from QCD instantons.

- Also, $10^9 \text{ GeV} < f_a < 10^{12} \text{ GeV}$.

2. Axions in String Theory

- String theory has lots of candidate axions.

  This is a nice feature

- Axions are components of moduli fields.

- In supergravity, a gauge kinetic function f gives couplings

$$\frac{\mathrm{Re}(f)}{4} \int F_{\mu\nu} F^{\mu\nu} + \frac{\mathrm{Im}(f)}{4} \int F_{\mu\nu} \tilde{F}^{\mu\nu}$$

- If Φ appears in f, $\mathrm{Im}(\Phi)$ has axionic couplings.

example

In the heterotic string, the dilaton superfield S is the universal gauge kinetic function

$$S = e^{-2\phi} \mathcal{V} + ia \qquad \left(\begin{array}{l} g_s = e^{\phi} \\ \mathcal{V} = \text{volume} \\ da = e^{-2\phi} * dB_{\mu\nu}, \end{array} \right)$$

The coupling

$$\frac{\text{Re}(S)}{4} \int F_{\mu\nu} F^{\mu\nu} + \frac{\text{Im}(S)}{4} \int \tilde{F}_{\mu\nu} \tilde{F}^{\mu\nu}$$

implies $\text{Im}(S) = a$ is an axion (the "model-independent axion").

There are also the "model-dependent axions". These arise from $\text{Im}(T_i)$, where $T_i = t_i + i b_i$ is a Kähler modulus.

example (IIB D3/D7 branes)

For D3-branes in type IIB string theory, the dilaton multiplet S is again the gauge kinetic function.

$$S = e^{-\phi} + i C_0$$

$\uparrow$ RR 0-form

The DBI and CS actions give

$$T_3 \int \sqrt{g} + T_3 \int C_4 + \underbrace{\frac{1}{4} \int e^{-\phi} F_{\mu\nu} F^{\mu\nu} \sqrt{g}}_{\mathrm{Re}(S)} + \underbrace{\frac{1}{4} \int C_0 F \wedge F}_{\mathrm{Im}(S)}$$

If QCD is realised on D3-branes, C_0 has the right couplings to be the QCD axion.

example

For D7-branes wrapping a 4-cycle Σ_i, an axion can come from $\mathrm{Im}(T_i)$.

Expanding the Chern-Simons coupling

$$\int e^{B+2\pi\alpha' F} \wedge \left(\sum C_i\right)$$

gives a term

$$\int F \wedge F \wedge C_4 .$$

Then $C_{4,i} = \int_{\Sigma_i} C_4$ has the axionic

coupling

$$C_{4,i} \int F \wedge F$$

$C_{4,i} = \mathrm{Im}(T_i)$, where T_i is a Kähler modulus.

String theory axions naturally have the Peccei-Quinn symmetry

$$a \to a + \epsilon .$$

This is exact in perturbation theory and is only violated by non-perturbative effects (e.g. worldsheet instantons / brane instantons).

Thus generic string compactifications do contain axions capable of solving the strong CP problem.

3. Axions and Moduli Stabilisation

- Axions belong to moduli multiplets.

- We want to
  - stabilise the moduli to avoid 5th forces
  - keep an axion unfixed to solve the strong CP problem.

- This creates a modulus 'anti-stabilisation' problem.

- Many approaches to stabilising moduli Φ_i stabilise $\mathrm{Re}(\Phi_i)$ and $\mathrm{Im}(\Phi_i)$ at the same time.

In IIB flux compactifications, the fluxes generate a potential for the dilaton S.

$$W = \int (F_3 + i S H_3) \wedge \Omega$$

The axion $\mathrm{Im}(S)$ gets a mass from the fluxes.

Conclusion 1

In IIB flux compactifications, QCD should not be realised on a D3-brane stack.

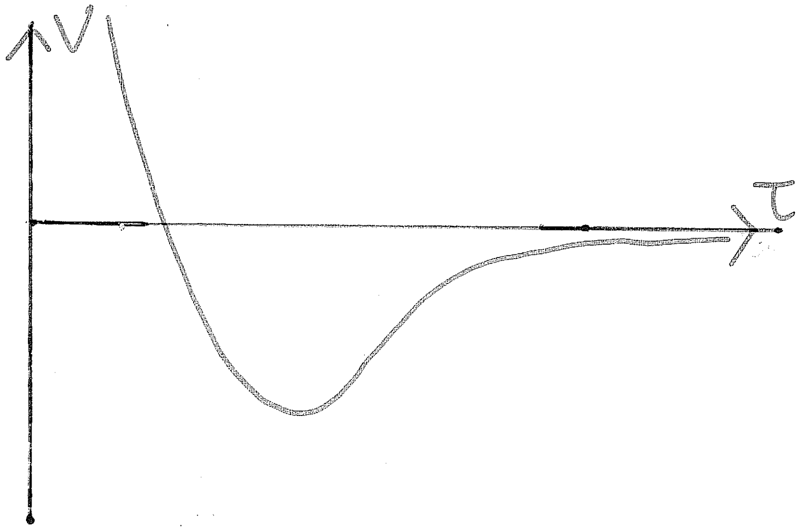
example 1-modulus KKLT

$$T = \tau + ic$$

$$W = W_0 + Ae^{-2\pi T}$$

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

· Solve $D_T W = 0$ to obtain a minimum of the potential.



· As T appears in the superpotential, a mass is generated for the axion $\mathrm{Im}(T)$. Therefore this axion cannot solve the strong CP problem.

Conclusion 2

T_{QCD} should not be stabilised by nonperturbative superpotential effects.

This suggests it should be stabilised by Kähler corrections.

A No-Go Theorem

- There exists no supersymmetric minimum of the F-term potential with stabilised moduli and unfixed axions.

Proof

① We have an $\mathcal{N}=1$ supergravity theory

$$W = W(\Phi_\alpha, T_\beta)$$
$$K = K(\Phi_\alpha, T_\beta + \bar{T}_\beta)$$

$b_\alpha = \mathrm{Im}(T_\alpha)$ are potential axions.

② We suppose we have

$$D_{\Phi_\alpha} W = 0, \quad D_{T_\beta} W = 0 \qquad \forall \alpha, \beta$$

but one axion $b_u = \sum_\beta \lambda_\beta b_\beta$ remains unfixed.

(3) At the susy locus,

$$V = -3e^{K}|W|^2$$

(recall $V = e^{K}(K^{-1} \bar{D}\bar{W} DW - 3|W|^2)$)

As K is independent of b_v, for b_v to be a flat direction $|W|$ must be independent of b_v.

(4) Either $W = e^{-\alpha T_v}$ (exceptional case)

or W is independent of b_v.

By holomorphy, W is then independent of $\tau_v = \mathrm{Re}(T_v)$ and T_v.

(5) Therefore, $\partial_{T_v} W \equiv 0$.

But $0 = D_{T_v} W = \cancel{\partial_{T_v} W}_{0} + (\partial_{T_v} K) W$

and so at the susy locus $\partial_{T_v} K = 0$.

($W=0$ is the exceptional case).

(6) We can now explicitly show

$$\frac{\partial^2 V}{\partial \tau_U^2} < 0$$

i.e. $\tau_U = \mathrm{Re}(T_U)$ is tachyonic.

Proof. $V = e^K \left(K^{i\bar{j}} D_i W D_{\bar{j}} \bar{W} - 3 |W|^2 \right)$

$$\Rightarrow \quad \frac{\partial V}{\partial \tau_U} = e^K K^{i\bar{j}} \left(\frac{\partial}{\partial \tau_U} (D_i W) D_{\bar{j}} \bar{W} + D_i W \frac{\partial}{\partial \tau_U} (D_{\bar{j}} \bar{W}) \right) - 3 \frac{\partial K}{\partial \tau_U} e^K |W|^2$$

(only keep terms non-vanishing at the susy locus)

$$\Rightarrow \quad \frac{\partial^2 V}{\partial \tau_U^2} = e^K K^{i\bar{j}} \left(2 \frac{\partial}{\partial \tau_U} (D_i W) \frac{\partial}{\partial \tau_U} (D_{\bar{j}} \bar{W}) \right) - 3 \frac{\partial^2 K}{\partial \tau_U^2} e^K |W|^2$$

Now use
- $D_i W = \partial_i W + (\partial_i K) W$
- $\frac{\partial}{\partial \tau_U} (W) = 0$
- $\frac{\partial}{\partial \tau_U} = 2 \frac{\partial}{\partial T_U}$

$$\Rightarrow \quad \frac{\partial^2 V}{\partial \tau_U^2} = 4 e^K \left(2 K^{i\bar{j}} K_{i\bar{U}} K_{U\bar{j}} - 3 K_{U\bar{U}} \right) |W|^2$$

$$= -4 K_{U\bar{U}} e^K |W|^2 < 0$$

Loopholes

- Try and use D-terms either to restore a minimum of the potential or to stabilise the moduli.

- The tachyon has mass

$$m^2_{\tau_0} = -2e^K |W|^2 = -\frac{8}{9} |m_{BF}|^2$$

and is Breitenlohner Freedman stable. However this stability criterion ceases to be relevant after uplifting.

Comments

- The result is pure supergravity and applies to all string compactifications.

- Unless the F-term potential is sensitive to all axions (i.e. through non-perturbative effects) it has no supersymmetric minimum.

- This has consequences for IIA flux compactifications – the supersymmetric solutions have many tachyonic directions.

Conclusion 3

To avoid tachyons, the moduli potential should break supersymmetry before the uplift.

Conclusion 4

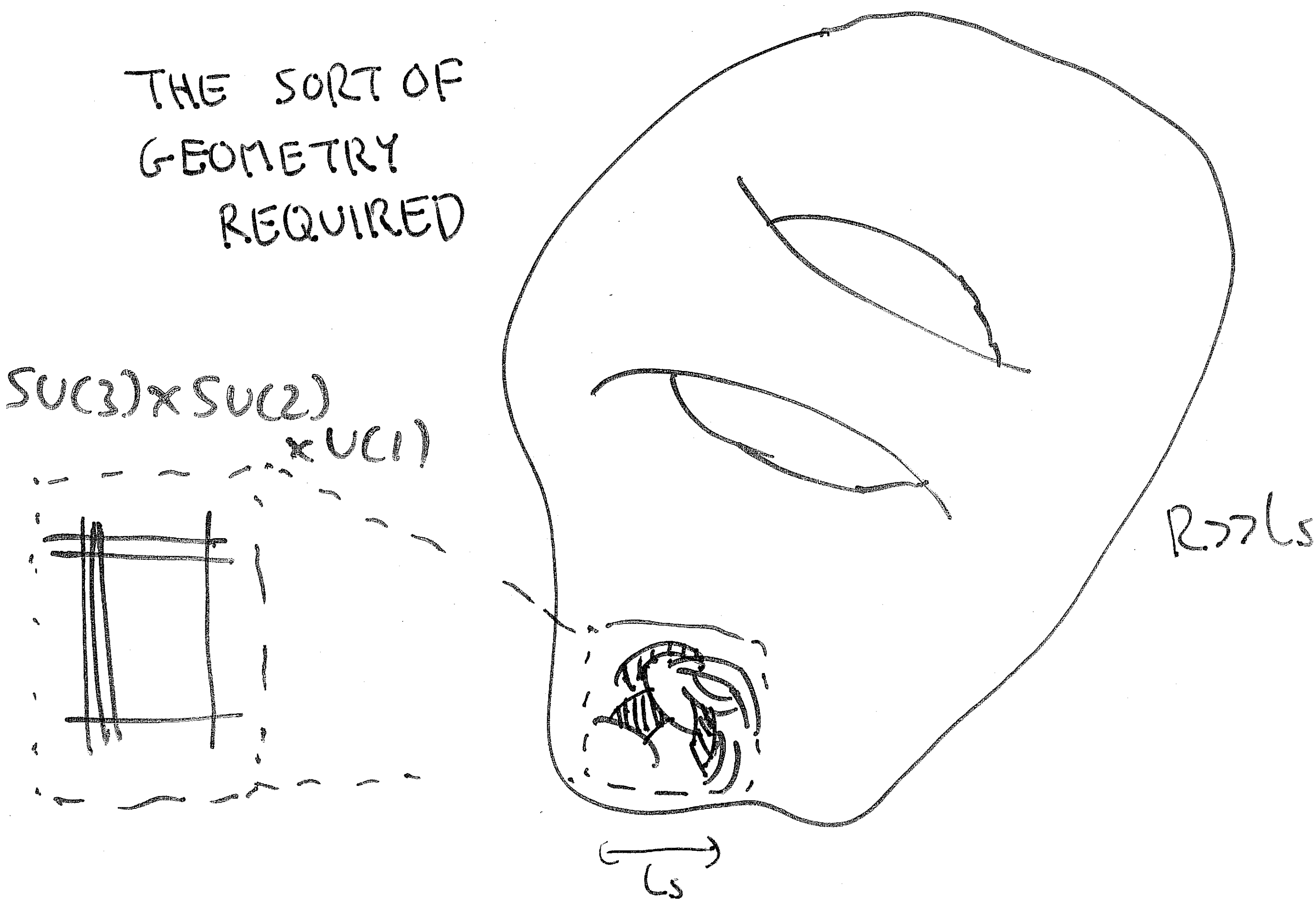
This favours gravity-mediated supersymmetry breaking.

The Axion Decay Constant

- f_a measures the axion-QCD coupling. This is essentially a stringy coupling. If we take the cosmological bound seriously, $10^9\,\text{GeV} < f_a < 10^{12}\,\text{GeV}$ suggests an intermediate string scale.

- It is hard to lower f_a if $M_s \sim M_P$.

THE SORT OF GEOMETRY REQUIRED



What strong CP tells us about the Landscape

1. In IIB flux compactifications, don't put QCD on D3 branes.

2. At least one modulus (T_{QCD}) should be stabilised perturbatively.

3. The moduli potential should break supersymmetry before uplifting.

4. Gravity-mediated supersymmetry breaking seems favoured.

5. A lowered string scale is preferred if we take the f_a bounds seriously.

III. From the Landscape to the LHC

- We want to predict the sparticle masses before they are measured.

- The structure of the moduli potential determines the breaking of supersymmetry.

- Different mechanisms of moduli stabilisation give different predictions for the soft masses.
  (c.f. Piyush's talk)

5. Review of Soft Masses

Given the supergravity scalar potential $V(\Phi_\alpha)$, the important quantities are the F-terms

$$F^\alpha = e^{K/2} \sum_\beta K^{\alpha\bar{\beta}} D_{\bar{\beta}} \bar{W}$$

Then gaugino masses are

$$M_a = \frac{1}{2\,\mathrm{Re}(f_a)} \sum_\alpha F^\alpha \partial_\alpha f_a$$

(f_a = gauge kinetic function)

and scalar masses

$$m_i^2 = m_{3/2}^2 - F^m \bar{F}^{\bar{n}} \partial_m \partial_{\bar{n}} (\ln \tilde{K})$$

($\tilde{K}$ = matter metric)

We focus on IIB D3/D7 flux compactifications. For a D7 brane stack,

$$f_a = \frac{T_a}{2\pi}$$

for branes wrapping cycle T_a.

Then

$$M_a = \frac{2\pi}{2\,\mathrm{Re}(T_a)} \sum_{\alpha} F^{\alpha} \partial_{\alpha} \left(\frac{T_a}{2\pi} \right)$$

$$= \frac{F^a}{2\tau_a}$$

The gaugino mass is determined by the F-term of the corresponding cycle.

In IIB models, the Kähler moduli are typically stabilised by non-perturbative effects.

$$K = -2\ln\left(\mathcal{V} + \frac{\xi}{2g_s^{3/2}}\right)$$

$$W = W_0 + \sum_i A_i e^{-a_i T_i}$$

We ask: if a modulus T_k is stabilised by non-perturbative effects; what does this tell us about F^k?

$$F^k = e^{K/2} \sum_j K^{k\bar{j}} D_{\bar{j}} \bar{W}$$

$$= e^{K/2} \sum_j K^{k\bar{j}} \left(\partial_{\bar{j}} \bar{W} + (\partial_{\bar{j}} K)\bar{W}\right)$$

Now, $K^{kj}\partial_j K = -2\tau_k$, and so we can write

$$F^k = e^{K/2}\Big(\underbrace{\sum_j -K^{kj} a_j A_j e^{-a_j T_j}}_{(K^{-1})^{kj}\partial_j W \quad (1)} \underbrace{-2\tau_k \overline{W}}_{(K^{-1})^{kj}(\partial_j K)\overline{W} \quad (2)}\Big)$$

Result

If T_k is stabilised non-perturbatively, there is a cancellation in the sum (1) + (2)

$$(1) + (2) = 2\tau_k \overline{W}\left(1 + O\left(\frac{1}{\ln(m_{3/2})}\right) - 1\right).$$

The gaugino mass is then suppressed

$$M_k = \frac{m_{3/2}}{\ln(m_{3/2})}$$

The suppression

$$M_k \sim \frac{m_{3/2}}{\ln(m_{3/2})}$$

is independent of whether the stabilisation is supersymmetric or not:
(approximately)

it follows purely from $\frac{\partial V}{\partial T_k} = 0$.

It was seen first in simple KKLT models but one can show
(CCNFP '04)
that the relation is generic.

In KKLT-style models (uplift a susy AdS solution) the scalar masses are also suppressed.

$$m_i^2 \sim \frac{m_{3/2}^2}{\ln(m_{3/2})^2}$$

This is not generic. In large-volume models

$$m_i^2 \sim m_{3/2}^2$$

Why? In KKLT, the anti-brane dominates the susy breaking. The F-term supergravity contribution is suppressed.

$$m_i^2 \sim m_{3/2}^2 - \underbrace{F^m F^{\bar{n}} \partial_m \partial_{\bar{n}} (\log \tilde{K}) - \frac{1}{3} V_{\bar{D3}}}_{\text{KKLT}}$$

7. Conclusions (again)

- Non-perturbative moduli stabilisation makes sense – it can naturally generate hierarchies (racetrack, IIB large volume)

- If T_k is stabilised non-perturbatively,

$$M_k \sim \frac{m_{3/2}}{\ln(m_{3/2})}$$

- To solve the strong CP problem, T_{QCD} should be stabilised perturbatively. This suggests the gluino mass should not be suppressed.

$$M_{\tilde{g}} \sim m_{3/2}.$$

- This motivates a scenario in which, at a high scale,

$$M_{\tilde{g}} \sim \ln(m_{3/2})\, M_{\tilde{W},\tilde{Z},\tilde{B}}$$

with heavy scalars $m_i \sim M_{\tilde{g}}$.

e.g.

$\tilde{q}, \tilde{L}, \tilde{g}$ $\quad \sim 3\,\text{TeV}$

$\chi_0, \chi_1, \tilde{W}$ $\quad \sim 100\,\text{GeV}$

- The Peccei-Quinn solution to the strong CP problem could be "seen" at the LHC – maybe.

- It certainly motivates "unnatural" sparticle spectra.

Summary

- The absence of CP violation in strong interactions may be a direct window on the landscape.

- The existence of a Peccei-Quinn axion tells us:
  - the moduli potential breaks susy
  - T_{QCD} is stabilised perturbatively

- Non-perturbative stabilisation gives suppressed gaugino masses $\frac{m_{3/2}}{\ln(m_{3/2})}$

- The gluino may be ~ 30 times heavier than the other gauginos.

Thanks to the organisers for a great workshop.