Gauge/Gravity Duality and Particle Physics

Shamit Kachru
KITP & UCSB

Based on work with:
Benini, Dymarsky, Franco, Simic, Verlinde: 0903.0619 & to appear;
Simic, Trivedi: 0905.2970;
Franco, to appear.

On the (n-th year of the) eve of LHC turn-on, it may be useful to ask: what do we expect to find, and can string theory possibly provide useful help for particle theorists?

In various frameworks, this question has actually been studied since the mid 1980s. I will not try to review the rich set of ideas that emerged over two decades. These include numerous constructions in the heterotic string, intersecting brane models, etc; and in each of these, different features of the physics are beautifully geometrized.

But to me, very broadly speaking, it seems that there are two circumstances where it is reasonable to try and apply string theory to particle phenomenology.

* The phenomenon under study is ultraviolet sensitive. This could include cases where phenomenology crucially depends on higher-dimension, Planck-suppressed operators (gravity mediation, inflation), or more subtle cases.

* The system under study is inherently strongly coupled. Then, quantitative calculations are difficult without importing new and powerful tools. Duality in field theory and string theory (notably, gauge/gravity duality) are two such tools.

Due to my parochial nature, I'll focus on two examples of the latter that I've thought about recently:

1. Non-supersymmetric models, with the hierarchy protected by strong coupling effects (which give large anomalous dimensions to dangerous operators that could ruin the hierarchy).

2. Supersymmetric models, which are even more direct in their mediation than direct gauge mediation ("single-sector"-like models)

I. Non-supersymmetric models that solve the hierarchy problem via strong coupling

This is in one sense an old idea, going back to theories of technicolor. Here, we want to explore a more recent idea, growing out of the work of Randall and Sundrum and Strassler.

Basic Idea: Consider a non-supersymmetric field theory governed by some fixed point Lagrangian at high energies:

$$\mathcal{L} = \mathcal{L}_{CFT}(\phi_i, \psi_n, A^a_\mu)$$

Suppose there is a finite list of relevant operators in the sense of the renormalisation group:

$$\mathcal{O}^1_{\Delta_1}, \mathcal{O}^2_{\Delta_2}, ..., \mathcal{O}^K_{\Delta_K} \qquad \Delta_i < 4$$

And suppose further that the field theory has a global symmetry group G, and these operators all transform in non-trivial representations.

Then, it would be reasonable to think that:

* Starting with this theory at a high energy scale Λ_{UV} , and preserving G symmetry, one has RG flow until some marginal operators which grow marginally relevant along the flow, become strong.

* These operators cause new behavior (confinement? flow to new different fixed point?) which could include generation of composite Higgs-like scalars, at exponentially small energy.

$$\Delta(\mathcal{O}) = 4 - \epsilon \ \rightarrow \ \Lambda_{IR} \sim \Lambda_{UV} \ e^{-\frac{1}{\epsilon}}$$

Problem: In simple non-supersymmetric theories with scalar fields in the UV, the operator

$$\mathcal{O} = \phi^\dagger \phi$$

is invariant under all (linearly realized) symmetries G, and is very relevant!

So instead of a long RG flow to generate an exponentially low scale before new phenomena occur, the RG flow "ends" right at the UV cutoff.

Observation of Strassler (2003):

In many known large N theories at large 't Hooft coupling, the global singlet operators

$$\mathcal{O} \sim \phi^\dagger \phi \;\rightarrow\; \Delta(\mathcal{O}) \sim (g^2_{YM} N)^{1/4}$$

The most obvious, dangerous global singlet is dual to a string state and becomes very irrelevant at large 't Hooft coupling!

This gives us a good reason to work at large 't Hooft coupling. Then, given a large N theory with no Global Singlet Relevant Operators (GSROs), we could make a field theory whose dual gravity description (via AdS/CFT) looks like:

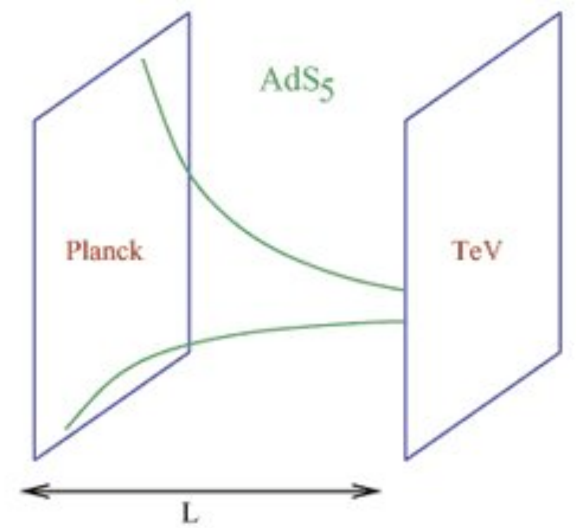
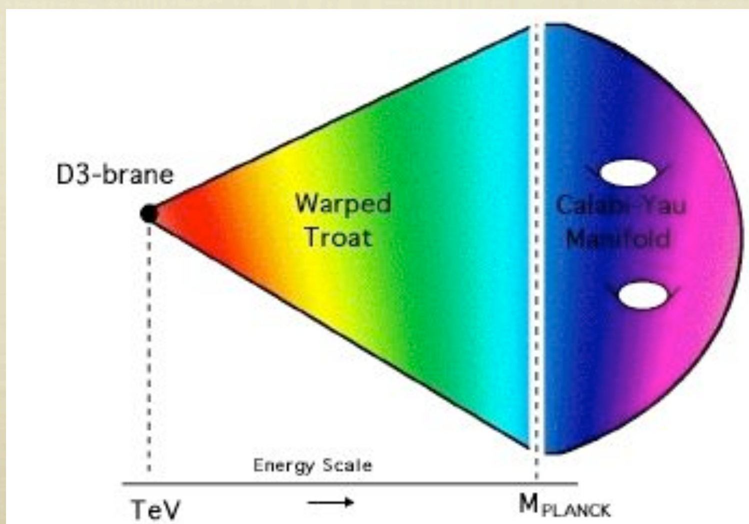


FIG. 8: Warped Extra Dimension. The local cutoff is exponentially smaller than M_P.

So at zeroth order, we need to find AdS/CFT dual pairs where:

1. There is a global symmetry G under which all relevant operators are charged.

2. We can preserve a sufficiently large subgroup of G when we "compactify the throat" that all relevant operators are still forbidden:



With Simic and Trivedi (0905.2970), we recently found:

1. A simple, infinite class of AdS/CFT pairs which have no global singlet relevant perturbations.

2. One case in which we can explicitly describe the IR physics, see ways of getting emergent scalars which Higgs a low-energy gauge group (though the structure is not at all Standard Model like), and compactify the theory preserving (a large enough subgroup of) the relevant symmetries.

A simple infinite class of such non-supersymmetric theories arises from **orbifolds of the N=4 field theory**.

Consider orbifolding the three complex planes transverse to N D3-branes by the group with generator:

$$\alpha = R_{(\frac{2\pi}{k})}(-1)^F .$$

The rotation acts explicitly as:

$$(Z^1, Z^2, Z^3) \to (\exp[\frac{2\pi i}{k}]Z^1, \exp[\frac{2\pi i}{k}]Z^2, \exp[\frac{2\pi i}{k}]Z^3).$$

* For k even, this orbifold actually takes the Type IIB string theory to a Type O string theory. We will not consider these cases.

* For k = 3, the resulting theory is an N=1 supersymmetric conformal field theory that has already been well studied.

* For k odd and k > 3, we obtain in this way candidate large N non-supersymmetric conformal field theories (at the planar level! But non-planar corrections are irrelevant for our purposes). We discuss only these cases.

This orbifold maps to a freely-acting orbifold on the near-horizon geometry (the fixed point at the origin is no longer present). The light states on the gravity side are then just the invariant KK modes in $AdS_5 \times S^5$.

Example: k=5

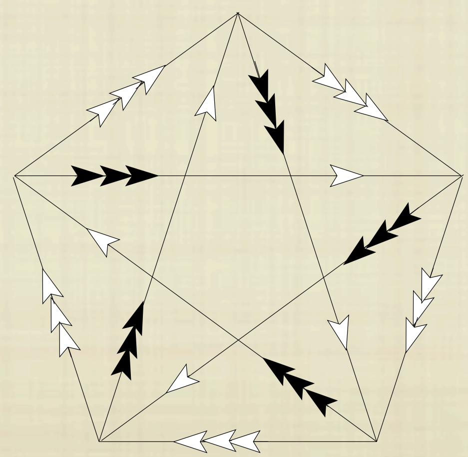


Figure 1: Quiver diagram of the $k = 5$ case. White arrows denote fermions, and black arrows denote scalars. We thank the authors of [13] for permission to reproduce this figure.

THIS THEORY HAS GLOBAL SYMMETRY GROUP

$$G = SU(3) \times U(1)$$

IT IS EASY TO ARGUE THAT THERE ARE NO GSROS.

FOR INSTANCE, JUST IN THE SCALAR BILINEAR SECTOR OF THE N=4 THEORY:

Operator	Δ	$SU(3) \times U(1)$
$\mathrm{Tr}(Z^i Z^j)$	2	$\mathbf{6}_{4/3}$
$\mathrm{Tr}(Z^i Z^{\bar{j}}) - \frac{1}{3}\mathrm{Tr}(Z^i Z^{\bar{i}})$	2	$\mathbf{8}_0$
$\mathrm{Tr}(Z^i Z^{\bar{i}})$	$(g_{YM}^2 N)^{1/4}$	$\mathbf{1}_0$

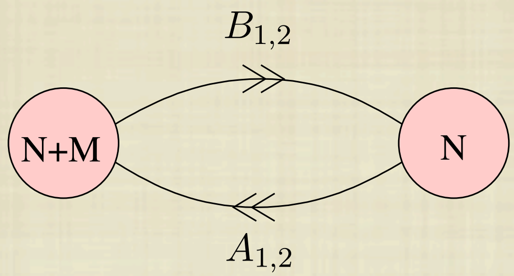
BRIEF INSPECTION SHOWS THAT ALL SCALAR BILINEARS IN THE ORBIFOLD THEORY INHERIT THESE SETS OF QUANTUM NUMBERS.

Similarly, one can see (by starting with the N=4 theory) that there are no GSROs in the fermion bilinear sector, the scalar trilinear sector, or the multi-trace operators.

This kind of argument shows that all of the odd k > 3 non-supersymmetric orbifolds satisfy our zeroth order criterion: they have scalars but no GSROs!

We don't yet have a good picture of what happens to these theories in the deep IR, when the (existing) singlet marginally relevant operators grow strong.

However, we can develop a picture also for the IR physics in examples based on the (deformed) conifold. The dual gauge theory is an N=1 susy theory with quiver:



$$W = \epsilon\epsilon ABAB$$

Field	$SU(2)\times SU(2)$	$U(1)_R$	$U(1)_B$
A_i	$(\mathbf{2},\mathbf{1})$	$1/2$	1
B_j	$(\mathbf{1},\mathbf{2})$	$1/2$	-1

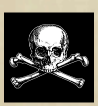
At finite M, anomalies lead to a breaking

$$U(1)_R \to Z_{2M}$$

There is also a symmetry which exchanges A,B accompanied by complex conjugation:

$$Z_{2,\text{exchange}}$$

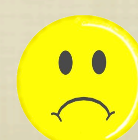
So, what is the situation with GSROs in the supersymmetric parent theory? Happily, Ceresole et al classified all operators in 1999 (via brute force dimensional reduction). Only a few low dimension SU(2) x SU(2) neutral operators exist:

	Operator	$U(1)_R$	$U(1)_B$
	$\text{Tr}(\vert A\vert^2 - \vert B\vert^2)$	0	0
	$\text{Tr}(\lambda\lambda)$	2	0
	$\text{Tr}(\epsilon_{ij}\epsilon_{kl}A_iB_kA_jB_l)$	2	0
	$\text{Tr}(AB)\overline{\text{Tr}(AB)}$	0	0

Two of those guys look dangerous:

$\mathrm{Tr}(|A|^2 - |B|^2)$ is protected - it is the lowest component of the baryon current supermultiplet. Fortunately, it is odd under $Z_{2,\mathrm{exchange}}$!

$\mathrm{Tr}(AB)\overline{\mathrm{Tr}(AB)}$ has dimension 3 at leading order in the 1/N expansion. It is a singlet under the entire group of global symmetries.

So the SUSY theory has GSROs. 

We will see, however, that obvious non-supersymmetric daughters do not.

Consider, for instance, the quotient of the conifold theory by:

$$A \to -A, B \to B \ .$$

$$\psi_A \to \psi_A, \psi_B \to -\psi_B \ .$$

$$\lambda \to -\lambda \ .$$

The resulting gauge theory (in the m=0 case) has the following spectrum.

2 x

	$SU(n)_1$	$SU(n)_2$	$SU(n)_3$	$SU(n)_4$
Q_1	$\square$	$\overline{\square}$	1	1
Q_2	1	$\square$	$\overline{\square}$	1
Q_3	1	1	$\square$	$\overline{\square}$
Q_4	$\overline{\square}$	1	1	$\square$

2 x

	$SU(n)_1$	$SU(n)_2$	$SU(n)_3$	$SU(n)_4$
ψ_1	$\overline{\square}$	$\square$	1	1
ψ_2	1	$\overline{\square}$	$\square$	1
ψ_3	1	1	$\overline{\square}$	$\square$
ψ_4	$\square$	1	1	$\overline{\square}$

	$SU(n)_1$	$SU(n)_2$	$SU(n)_3$	$SU(n)_4$
λ_1	$\square$	1	$\overline{\square}$	1
λ_2	$\overline{\square}$	1	$\square$	1
λ_3	1	$\square$	1	$\overline{\square}$
λ_4	1	$\overline{\square}$	1	$\square$

All of the symmetries of the conifold theory descend to this theory. In the case M=0, the only change is that two of the gauge groups become SU(N+M), and the anomaly now breaks:

$$U(1)_R \to Z_M$$

* Descendants of operators we discussed before are projected out exactly as before.

* What about the fearful $\mathrm{Tr}(AB)\overline{\mathrm{Tr}(AB)}$?

The single trace operator $\mathrm{Tr}(AB)$ in the parent theory, is not invariant under the orbifold group! As a result, there is no gauge invariant operator in the daughter theory which is bilinear with one descendant of the A and one descendant of the B fields. To get a gauge invariant single trace operator then involves a string of 4 As and Bs (and yields no GSRO); the double-trace global singlet operator one can make from this is then highly irrelevant!

Simple orbifolds of the deformed conifold theory, then yield AdS/CFT duals which have all of the properties we require:

1) The field theory has no GSROs.

2) The theory has a marginal operator which grows marginally relevant, and smoothly ends the geometry in a tip which is well described by supergravity.

3) One can embed the dual throat geometry in an (orbifold of a) compact Calabi-Yau space, preserving large enough subgroups of G to forbid GSROs.

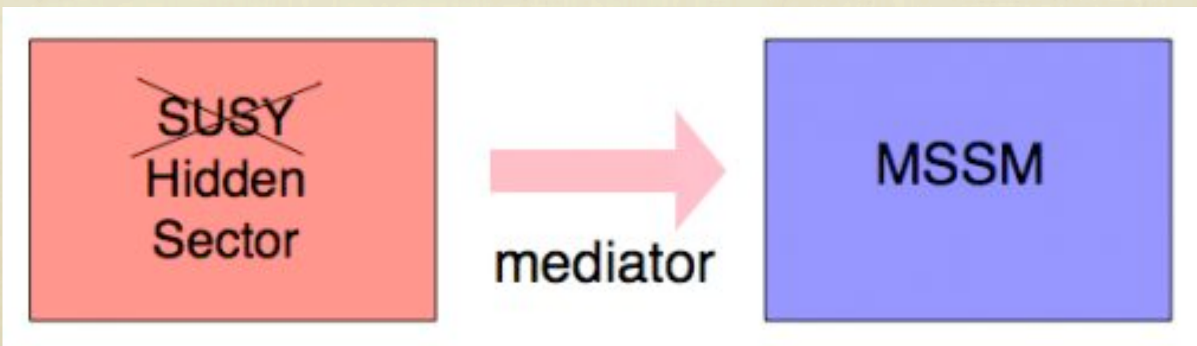
4) With a bit more work, one can show that one can get an IR field theory with emergent scalars and gauge fields, with the Higgs phenomenon occurring at the IR scale.

It is far from clear that one can make realistic models in this framework. One would likely want to add bulk gauge fields (flavor branes) to the throat, and localized fermions, all without destroying the protection from GSROs.

But as such models would be a completely different solution to the hierarchy problem than SUSY, that would also work up to very high energy scales (not just a decade in energy), they seem like an interesting idea to pursue.

II. Supersymmetric models with composite quarks and leptons

The existing paradigm of supersymmetric model building always starts with a picture of the form:



Are there ways to build SUSY models which appear less modular? While such modularity does not appear unnatural in top-down constructions, it would be perhaps more elegant to have a more closely integrated structure, where the SM fields themselves are "closer to" the SUSY breaking dynamics.

Dimopoulos and Georgi argued that if:

1) The gauge group is that of the SM
2) No higher dimension operators appear in K
3) Tree approximation is reliable

then there will always be a colored scalar sparticle ligher than the down quark.

* Gravity mediation violates assumption 2)

* Gauge mediation violates assumption 3)

Could there be other ways to circumvent their theorem, which involve more closely integrated SM and SUSY breaking dynamics?

Why not make the observable fields composite? As we'll see, this would also have the advantage of (partially) explaining mysterious facts about flavor physics.

examples (Non-Calculable) were provided by Arkani-Hamed, Luty and terning (1998), and Luty and Terning (1999).

* SUSY broken by strong dynamics at scale Λ

* SUSY breaking theory has an unbroken global symmetry group G

* G has cubic anomalies, and one has composite fermions in the IR via 't Hooft anomaly matching

* The SM gauge group arises as a subgroup of G which is anomaly free (or whose anomaly is cancelled by "elementary" states)

Of course, it would be nice if there were simple, calculable examples of this phenomenon.

Two approaches to making such models:

* Work directly in supersymmetric field theory, using techniques of holomorphy and duality. This seems to work very well!

S. Franco,
SK to appear

* Use gauge/gravity duality to geometrize the strong coupling dynamics, and build composite models in warped geometries.

↗ we'll talk about this...

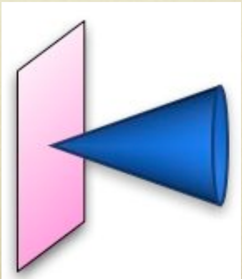
Gabella, Gherghetta, Giedt (2007);
Benini, Dymarsky, Franco,
SK, Simic, Verlinde, 0903.0619
and to appear

To start with, we need a strongly coupled hidden sector that has a low-scale SUSY-breaking state.

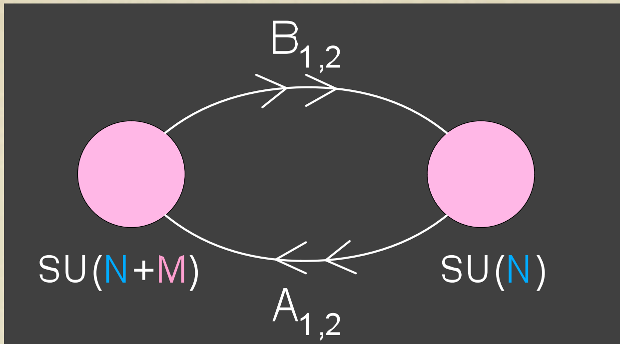
A good string toy model for this arises in the deformed conifold geometry, again. Lets go into more detail this time:

$$x^2 + y^2 + z^2 + w^2 = 0$$

This is the (undeformed) conifold. Now, add N D3 branes at the tip, and wrap M D5 branes on the small sphere at the tip:



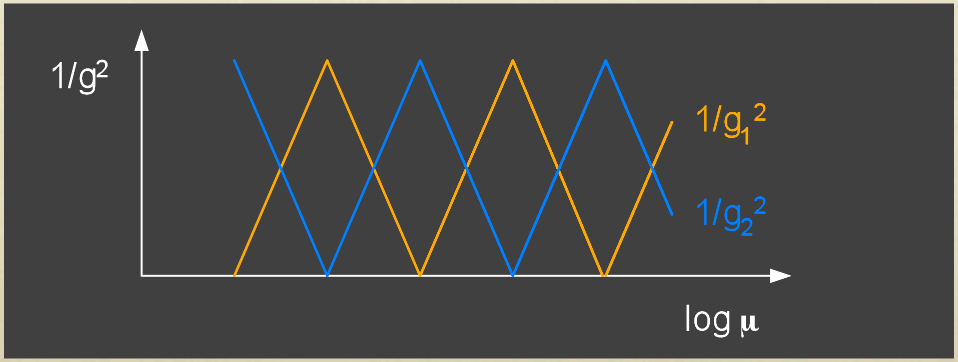
The dual gauge theory is:



$$W = \epsilon^{ij}\epsilon^{kl} A_i B_k A_j B_l$$

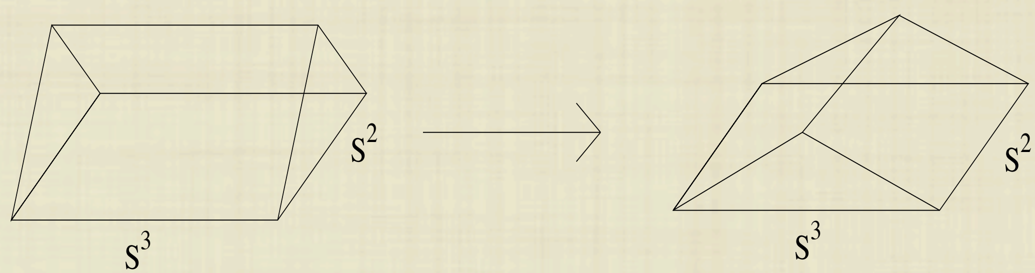
Klebanov, Witten;
Klebanov, Strassler

This theory undergoes a cascade of Seiberg dualities, as one gauge factor and then the other becomes strongly coupled:



The rank of the groups successively shifts down by M units, while leaving a self-similar superpotential. For N = KM, this continues (on one branch) until the IR dynamics confines and produces a gap. The gravity dual of the IR is a geometric transition:

(Pure geometry with fluxes)



(D-brane side)

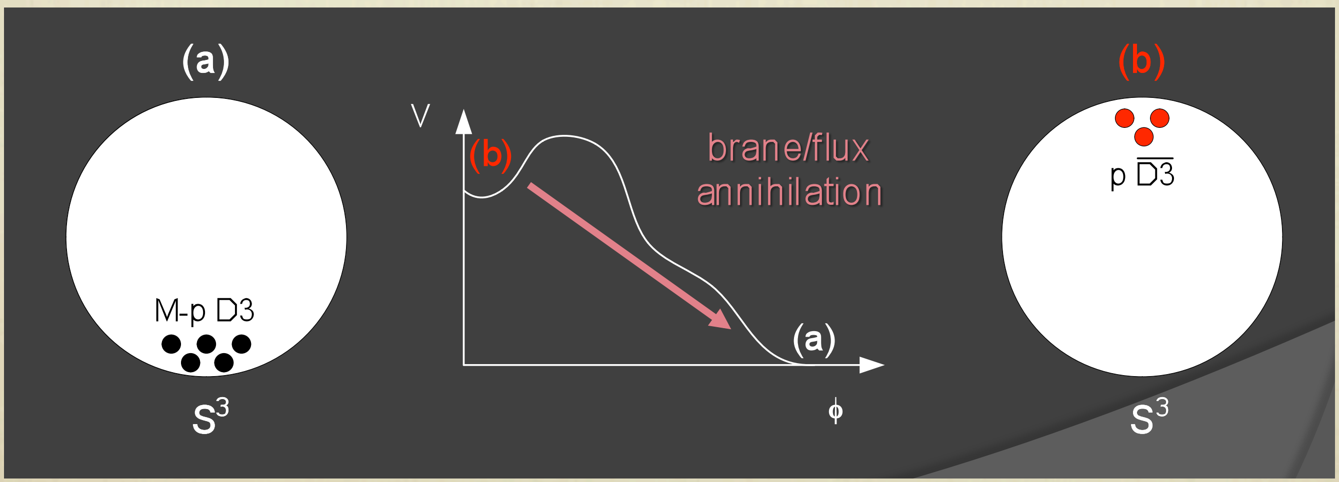
Vafa; Klebanov, Strassler

$$\int_A F_3^{RR} = M \qquad \int_B H_3^{NS} = -K$$

THE GEOMETRY BECOMES:

$$x^2 + y^2 + z^2 + w^2 = \epsilon^2, \quad \epsilon \sim \mathrm{Exp}(-4\pi K/3g_s M)$$

IF INSTEAD N = KM - P (P << M,N), ON ONE BRANCH WE ARE LEFT WITH (M-P) PROBE D3-BRANES AT THE TIP. THERE IS A METASTABLE SUSY-BREAKING STATE IN THIS THEORY:



SK, PEARSON, VERLINDE

In the limit of large distance from the tip, the metastable state is characterized by normalizable perturbations to the supersymmetric background:

DeWolfe, SK, Mulligan

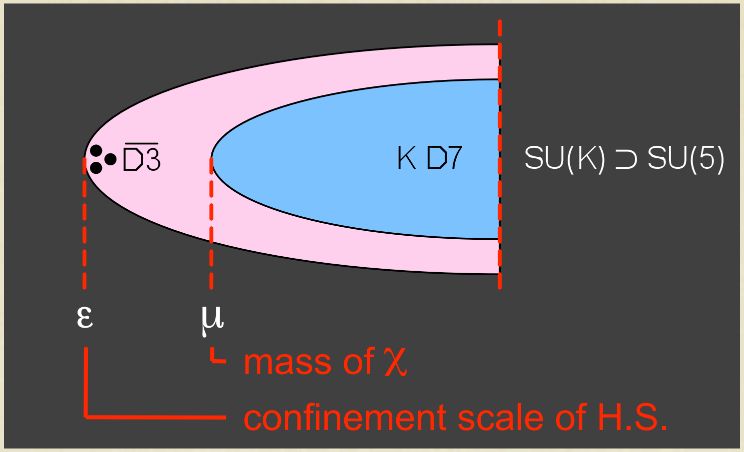
$$ds^2 = r^2 e^{2a(r)} \eta_{\mu\nu} dx^\mu dx^\nu + e^{-2a(r)} \left(\frac{dr^2}{r} + \sum_{i=1}^{2} (e_{\theta_i}^2 + e_{\phi_i}^2) + e^{2b(r)} e_\psi^2 \right)$$

$$\begin{aligned}
e^{-4a} &= \frac{1}{4} g_s \bar{N} + \frac{1}{8}(g_s \bar{M})^2 + \frac{1}{2}(g_s \bar{M})^2 \log r \\
&\quad + \frac{1}{r^4} \left[\left(\frac{1}{32} g_s \bar{N} + \frac{13}{64}(g_s \bar{M})^2 + \frac{1}{4}(g_s \bar{M})^2 \log r \right) \mathcal{S} - \frac{1}{16}(g_s \bar{M})^2 \phi \right] \\
e^{2b} &= 1 + \frac{1}{r^4} \mathcal{S} \\
k &= g_s \bar{M} \log r + \frac{1}{r^4} \left[\left(\frac{3}{8} \frac{\bar{N}}{\bar{M}} + \frac{11}{16} g_s \bar{M} + \frac{3}{2} g_s \bar{M} \log r \right) \mathcal{S} - \frac{1}{4} g_s \bar{M} \phi \right] \\
\Phi &= \log g_s + \frac{1}{r^4} [\phi - 3\mathcal{S} \log r]
\end{aligned}$$

$$\mathcal{S} \sim \frac{p}{N} e^{\left(-\frac{8\pi N}{3 g_s M^2} \right)}$$

Vacuum Energy

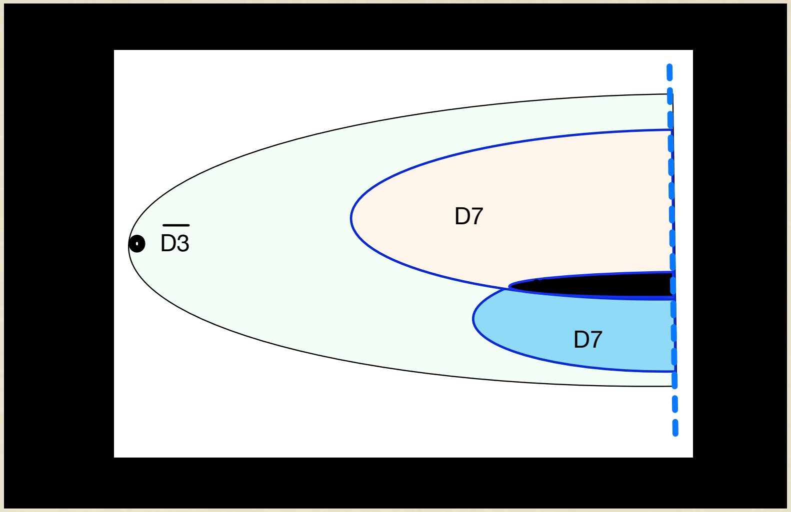
On top of this SUSY-breaking background, we can build a theory which incorporates a toy composite Standard Model in the following way.



We add (supersymmetric) Ouyang-embedded D7-branes:

$$z_1 z_2 - z_3 z_4 = \epsilon^2, \quad z_4 = \mu$$

And to get the "right" charged matter, e.g. 5s of SU(5), we intersect them with "flavor" D7s:



$$\text{flavor } D7 \;:\; z_1 = \tilde{\mu}$$

The number of chiral matter fields localized on the intersection curve C is:

$$n = \int_C (\mathcal{F}_{\rm SM} - \mathcal{F}_{\rm flavor})$$

There is a supersymmetric flux:

$$P \text{ of hodge type } (1,1), \quad P \wedge J = 0$$

Chen, Ouyang, Shiu

that we can place on the flavor brane, and that reduces on the intersection curve to:

$$\frac{P}{2i\tilde{\mu}^2} = \left(\frac{3}{2} + \frac{|z_3|^2}{\tilde{\mu}^2}\right) \frac{dz_3 \wedge d\bar{z}_3}{(|z_3|^2 + \tilde{\mu}^2)^2}$$

For a given radial (UV) cut-off, by scaling P appropriately we can localize any number of chiral zero modes.

We'll argue now that the degree of compositeness of matter added in this way, depends on the minimal radial location of the intersection.

It will be easier for these purposes to truncate to a toy 5D model, to avoid overly complicated formulae.

5D Toy Models

The basic strategy to make a toy model of our setup, is to take a slice of AdS with IR boundary conditions that break SUSY. The matter fields are in bulk 5D N=1 SUSY multiplets. For fermions, these are labelled by the bulk mass m (or equivalently c, the mass in units of the AdS curvature scale):

Gherghetta, Pomarol

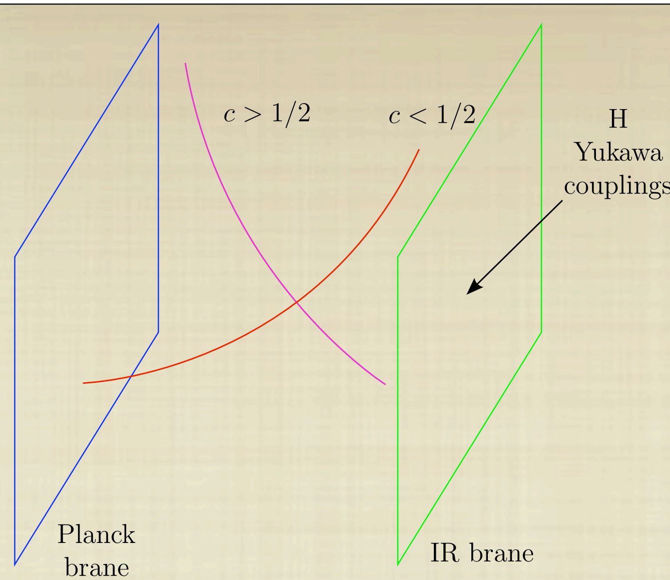


Figure 20: Fermion zero-mode profiles for different 5D fermion masses

THIS PICTURE (SUNDRUM TASI LECTURES) SHOWS HOW, DEPENDING ON "C", THE MATTER FIELDS ARE EITHER IR OR UV LOCALIZED. IN A MODEL WHERE THE HIGGS IS ALSO LOCALIZED, LARGE YUKAWAS OCCUR FOR MODES LOCALIZED NEAR THE HIGGS, AND SMALL YUKAWAS FOR THOSE WHICH ARE DISTANT. (THE PICTURE IS OF AN RS SCENARIO WHERE H IS IR LOCALIZED).

Quantitatively, in the warped metric:

$$\begin{aligned} ds^2 &= e^{-2kR\phi}\eta_{\mu\nu}dx^\mu dx^\nu - R^2 d\phi^2, \quad 0 \le \phi \le \pi \\ &\underset{y \equiv R\phi}{=} e^{-2ky}\eta_{\mu\nu}dx^\mu dx^\nu - dy^2, \quad 0 \le y \le \pi R\,. \end{aligned}$$

these wavefunctions take the form:

$$\hat{\Psi}_{\mathrm{L}}(x,y) = \hat{\Psi}_{\mathrm{L}}^{(0)}(x)\, e^{(\frac{1}{2}-c)k|y|} \quad ,$$

explaining the sharp localization as c moves from less than to more than 1/2.

The fermions with small c (and their bosonic partners) are naturally called composite. This has a natural explanation via the AdS/CFT correspondence. Consider a fermion with a given value of c. It is dual to an operator in the dual CFT, with conformal dimension:

$$\dim(\mathcal{O}) = \frac{3}{2} + \left|c + \frac{1}{2}\right|$$

Therefore if we start by assuming the 4D fermion to have canonical dimension, the boundary coupling:

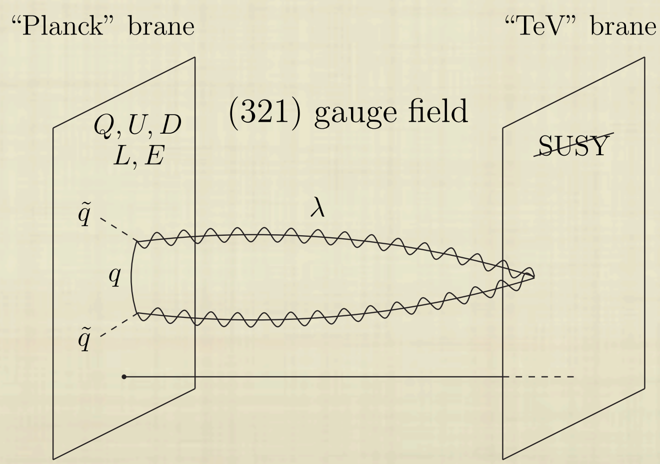
$$\mathcal{L} = ... - \Psi\mathcal{O} + ...$$

is RELEVANT for c < 1/2. The fermion mixes in an important way with CFT d.o.f.

For us it will be natural to take the Higgs to be **UV localized**; its small mass is explained by the small SUSY breaking scale, which has been geometrized via warping.

Then one can easily imagine two scenarios.

1. Holographic gauge mediation



c.f. Nomura 04

Benini, Dymarsky, Franco, SK, Simic, Verlinde

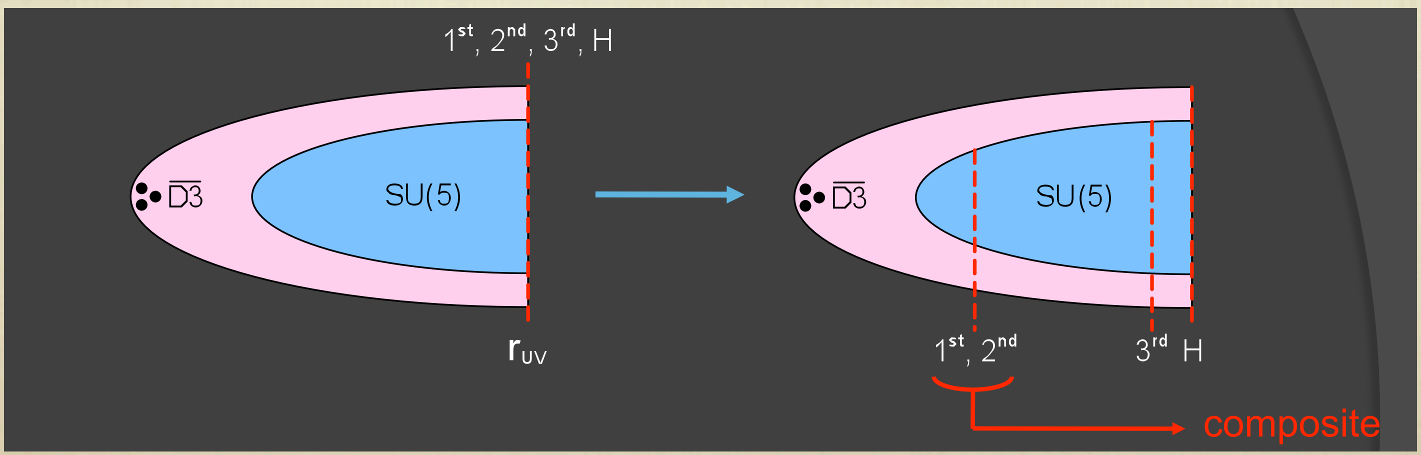
all matter fields UV localized; the bulk gaugino feels and transmits SUSY breaking.

2. Composite Models

Here, we spread the matter fields around in the microscopic 10D theory, by choosing different flavor-brane embeddings:

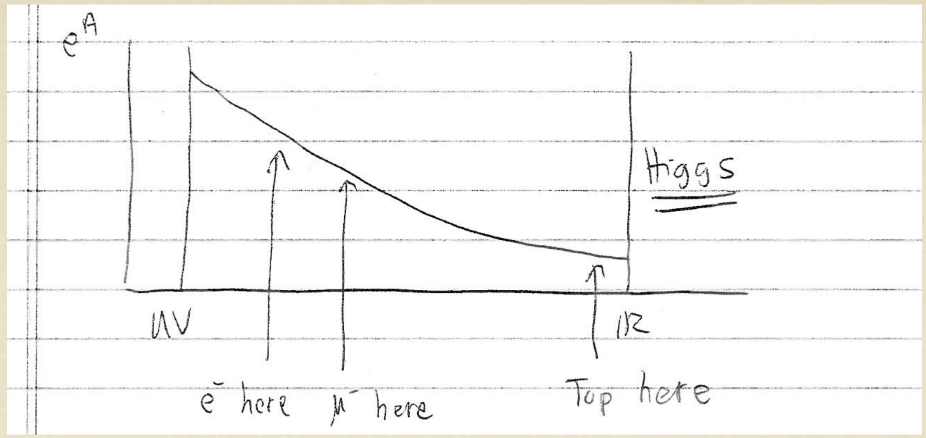
$$z_1^i = \tilde{\mu}^i, \quad \tilde{\mu}^3 >> \tilde{\mu}^{1,2}$$

In terms of pictures, the two scenarios then look like:

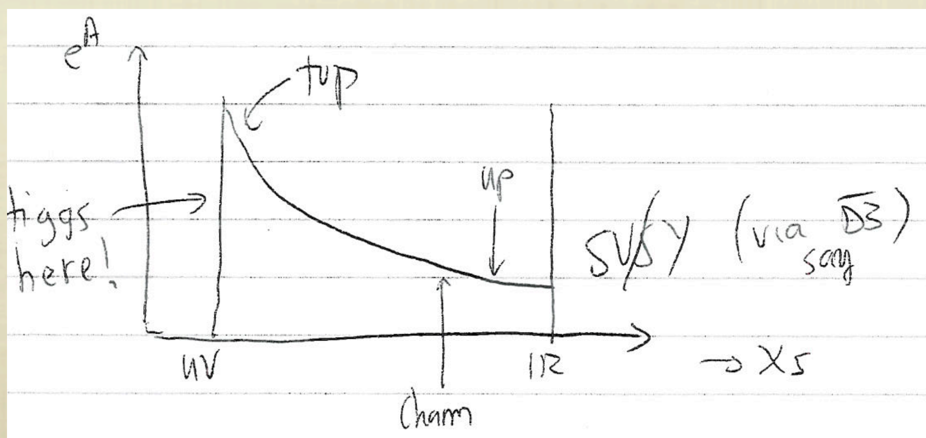


Note that in the SUSY composite models, one imagines **opposite** higher-dimensional geography, as compared to the RS models:

RS →



Single Sector →



Let us momentarily consider a 5D truncation of a non-supersymmetric AdS solution of string theory, following Gabella, Gherghetta and Giedt.

The AdS is deformed to a slightly different geometry, by the SUSY-breaking perturbation:

$$ds^2 = A^2(z)\left(-dt^2 + d\mathbf{x}^2 + dz^2\right)$$

$$A^2(z) = \frac{1}{(kz)^2}\left(1 - \epsilon(\frac{z}{z_1})^4\right)$$

$$z_0 \leq z \leq z_1$$

(intuitively, z_0 is UV and z_1 is IR).

This perturbation to the metric backreacts on the 5D fermionic and bosonic wavefunctions.

Now, the bosonic wavefunctions are characterized by a power-law behavior

$$\phi(z) \sim z^{b-1}$$

$$b = \frac{3}{2} - c$$

i.e., partners of UV localized fermions ($c > 1/2$) are also UV localized, while partners of IR localized fermions ($c < 1/2$) are also IR localized.

The SUSY-breaking perturbation of the 5D warp factor, induces a shift to the boson masses. For the UV localized modes, it is a very small effect:

$$\Delta m = \sqrt{\epsilon(1-b)(b+10)}(kz_1)^{b-1}z_1^{-1}$$

But for the IR-localized bosons, the mass shift is of order the scale of IR SUSY breaking:

$$\Delta m = \sqrt{\epsilon(b-1)(b+10)}z_1^{-1}$$

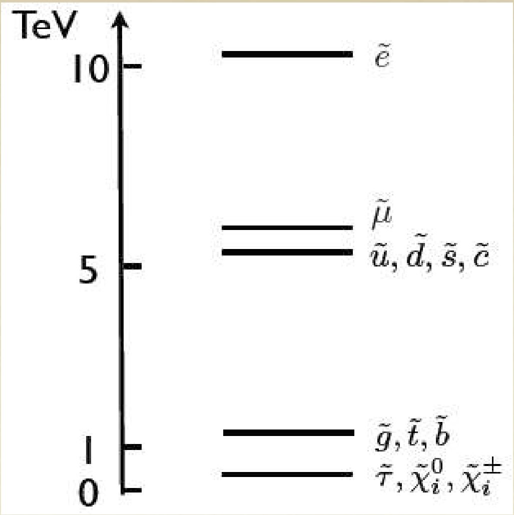
RESULT: in the models where the first two generations are composite, the spectrum is quite distinctive.

* The Stop gets its mass through gaugino mediation, and is the lightest matter sparticle.

* The first two generation sparticles are very heavy, receiving large “compositeness contributions” to their masses, in typical models.

* The Gauginos are much lighter than the first two generation sparticles.

So there is a direct correlation: small Yukawas -> large sparticle masses!



(From GGG)

FIG. 1: The generic mass spectrum of the 5D gravity model showing the heavy first and second generation scalars and lighter third generation scalars, gluinos, neutralinos and charginos. The LSP is the gravitino (not shown).

Note that it is an old idea (Dimopoulos and Giudice 95; Cohen, Kaplan, Nelson 96) that if the squarks in the first two generations are quite heavy, flavor problems can be vitiated.

So why not push up the first two generation masses even more (make them “more composite”)?

basic problem (Arkani-hamed, Murayama 97):

to be heavy enough to sufficiently suppress FCNCs, the first two generation sparticles must have masses > 22 TeV. But a two-loop contribution to the stop mass running then makes the stop tachyonic unless it is heavier than 4 TeV --> tuning in the Higgs sector.

THIS PROBLEM IS MODEL-DEPENDENT, AND THERE ARE WAYS AROUND IT BY CLEVER MODEL BUILDING.

Hisano, Kurosawa, Nomura (1999)

But the upshot is that we probably want to limit the size of the first two generation sparticle masses, and impose a flavor symmetry.

In any case, there is much work to do to get these models on solid footing like their gauge-mediated cousins; and string theory provides a useful tool for calculating in these intrinsically strongly-coupled systems!