

Dynamical SUSY Breaking in Intersecting Brane Models

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Abstract

We present a simple mechanism by which supersymmetry can be dynamically broken in intersecting brane models, naturally generating an exponentially small scale. Rather than utilize either non-Abelian gauge dynamics or D-instantons, our mechanism uses worldsheet instantons to generate the small scale in a hidden sector.

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1 Introduction

For many years, there has been great interest in dynamical supersymmetry breaking (DSB) as a method of solving the hierarchy problem by generating the electroweak scale[1]. There have been many realizations of this mechanism in field theory, as well as attempts to embed this method in string theory models. One very common feature of known DSB examples is the presence of a non-Abelian gauge group. It is the RG flow of this group which generates the low-scale dynamically via dimensional transmutation.

One example of this type of DSB is the model of ISS, and the several subsequent related models[2]. One is tempted to try to embed this type of model in the hidden sector of a string construction and thus obtain an example of DSB in string theory. For example, one could imagine constructing an intersecting brane model[3, 4] (IBM) with a Standard Model visible sector and a hidden sector brane whose low-energy effective theory is SQCD in the window $N_c + 1 \leq N_f \leq \frac{3}{2}N_c$. Unfortunately, this is not a trivial thing to do in known IBM constructions. Almost all explicit IBM's are constructed on relatively simple toroidal orientifolds, in which the orientifold planes generate negative space-filling charges of $\mathcal{O}(10)$. Because these charges are cancelled by the presence of Standard Model sector and hidden sector branes, they bound (usually severely) the number of colors one can arrange in the hidden sector. On the other hand, one often gets many flavors in the hidden sector due to the moderately large topological intersection number between different hidden sector branes. It turns out to be non-trivial to obtain a specific Type IIA intersecting brane model which manifests DSB of the form discussed by ISS, et al. (we do not know of an example). Although this is not expected to be problematic for IBM's constructed on more complicated manifolds, there has been very limited work in this area[5]. As such, it would be very nice to have a model of DSB which arises in a sector with only $U(1)$ gauge groups, which are plentiful in the IBMs which are easiest to construct explicitly.

Another difficulty with the class of models discussed in [2] is that one generally has multiple energy scales which must be generated, and with a particular hierarchy between them (essentially, the mass m of the hidden sector quarks and the dynamical scale Λ of that sector, with $m \ll \Lambda$). Generating these multiple scales typically complicates the model, and obtaining the appropriate hierarchy of scales usually requires an additional tuning and/or further non-perturbative dynamics. A DSB mechanism which depended only on one new scale would potentially be both more elegant and more easily realized in specific models.

In an interesting recent paper[6], Aharony, Kachru and Silverstein pointed out that standard field theory models of supersymmetry breaking (the Fayet, Polonyi and O'Raifeartaigh

models) can be realized dynamically in string theory constructions which do not rely on non-Abelian dynamics to generate the dynamical scale. In the cases they discussed, the gauge theory is realized by branes at singularities and D-brane instanton effects generate an exponentially small scale.

In this brief note, we discuss a similar method for generating a small SUSY-breaking scale dynamically, which appears quite naturally in intersecting brane models. In this case, worldsheet instantons generate the small dynamical scale. In section 2 we describe the intersecting brane model setup, and in section 3 we show how it naturally leads to dynamical supersymmetry breaking. We conclude with a discussion of the implications for intersecting brane models and phenomenology in section 4.

2 IBM Setup

The basic idea of an intersecting brane model (in Type IIA), is to compactify 10D Type IIA string theory on an orientifolded Calabi-Yau 3-fold. Spacetime-filling D6-branes must be added to cancel the charges of the orientifold planes, and the gauge and matter dynamics of these branes are relied upon to yield a visible SM-like sector, plus various hidden sectors. Importantly, the chiral matter content of the theory is counted by the topological intersection numbers of the branes; at each topological intersection point of any two branes (or their orientifold images), there lives an $N = 1$ chiral multiplet transforming in the bifundamental of the gauge groups living on the two branes. Furthermore, any two spacetime-filling D6-branes will generically have non-zero intersection number, since they wrap 3-cycles on a 6-manifold. This chiral matter can lead to mixed anomalies, which are cancelled by the Green-Schwarz mechanism and can give the gauge boson a mass. However, cancellation of the RR-tadpoles (or, equivalently, Gauss' Law) implies that there are no cubic anomalies. A similar setup exists for brane models in Type IIB.

We consider the simple case of three hidden sector branes, a , b and c with gauge groups $U(1)_a$, $U(1)_b$ and $U(1)_c$ and gauge couplings $g_{a,b,c}$ respectively. Generically, they all intersect each other, and we will assume they have intersection numbers $I_{ab} = I_{bc} = I_{ca} = 1$ (these numbers are chosen for simplicity; they are not essential for the argument). The D -term potential is then given by

$$V_D = \frac{g_a^2}{2}(|\phi_1|^2 - |\phi_2|^2 - \xi_a)^2 + \frac{g_b^2}{2}(|\phi_2|^2 - |\phi_3|^2 - \xi_b)^2 + \frac{g_c^2}{2}(|\phi_3|^2 - |\phi_1|^2 - \xi_c)^2 \quad (1)$$

where $\phi_{1,2,3}$ are the scalars of the chiral multiplets living at the three intersections, and $\xi_{a,b,c}$ are the FI-terms of the various $U(1)$'s[7]. They are constrained by $\sum \xi = 0$. Furthermore,

gauge-invariance implies that one cannot write a tree-level mass term in the superpotential. Instead, the first renormalizable superpotential term which one can write is

$$W = \lambda\phi_1\phi_2\phi_3. \quad (2)$$

Here λ is generated by a world-sheet instanton stretching between the three branes a , b and c . In particular, λ is expected to scale as $e^{-\frac{A}{\alpha'}}$ where A is the area of the instanton[4, 8]. This exponential suppression (in the large volume limit) implies that generally $\lambda \ll g_{a,b,c}$. Further Kähler corrections to the F -term potential will be suppressed by powers of M_{pl} .

Since the matter content is non-vectorlike, the non-diagonal $U(1)$ factors will have mixed anomalies which are cancelled by the Green-Schwarz mechanism. As such, there will be closed string axions which shift under the $U(1)$ gauge symmetries to restore gauge invariance. One should worry that these axions could appear in exponent of the superpotential couplings to terms which are naively gauge-noninvariant[11]:

$$\begin{aligned} W_1 &\sim \sum_i [e^{a_i}] \phi_i \\ W_2 &\sim \sum_{i,j} [e^{a_k}] \phi_i \phi_j. \end{aligned} \quad (3)$$

In this case, the exponentials in brackets involve sets of axions whose shifts under gauge transformations compensate for the phase which the scalars get. These terms are generated by non-perturbative instanton couplings, and thus are also exponentially suppressed.

There are regions of closed string moduli space where these linear and quadratic couplings can be more highly suppressed than Yukawa couplings. Either ϕ_k or $\phi_i\phi_j$ transforms under two $U(1)$ gauge groups which live on, say, branes i and j . In order to restore gauge invariance, the coefficient must contain two axionic couplings, and holomorphy implies that each coupling is accompanied by a suppression exponential in a Kähler modulus in Type IIB, or a complex structure modulus in Type IIA. Although it is possible in some cases to suppress these linear and quadratic couplings relative to the Yukawa, it may not always be possible and is in any case unnecessary to our main point. The only thing to note is that these couplings are still exponentially small compared to the string scale, and so we will assume for simplicity that the largest of these small couplings is about the same order as λ .

Note that we have not addressed the issue of closed string moduli stabilization. In general, the FI-terms will depend on the closed string moduli (complex structure moduli in Type IIA, Kähler moduli in Type IIB). It is important for these moduli to be stabilized, in order to prevent a runaway in closed string moduli space which could effectively set the D -term potential to zero and restore supersymmetry[9]. Stabilization of these moduli is in any case

required for phenomenological reasons. There are several known methods for stabilizing these moduli in Type IIA/B and related contexts[10].

3 Dynamical Breaking

If we had $\lambda = 0$, then we would have $V_F = 0$. There exists a solution (indeed a one-parameter family of them) for which $V_D = 0$, and these would correspond to supersymmetric vacua. Let us assume, without loss of generality, $\xi_a > 0$, $\xi_{b,c} < 0$. Expanding around $\lambda = 0$, we see that one minimum of the potential will be

$$\begin{aligned} |\phi_1|^2 &= \xi_a + \mathcal{O}\left(\frac{\lambda^2 \xi}{g^2}\right) \\ |\phi_2|^2 &= \mathcal{O}\left(\frac{\lambda^2 \xi}{g^2}\right) \\ |\phi_3|^2 &= -\xi_b + \mathcal{O}\left(\frac{\lambda^2 \xi}{g^2}\right) \end{aligned} \tag{4}$$

where we have taken the $\xi_{a,b,c} \sim \xi$ to be about the same scale and the $g_{a,b,c} \sim g$ to be about the same order. There is a compact flat direction corresponding to the overall complex phase of the scalars. Furthermore, there is a D -flat direction corresponding to increasing $|\phi_1|^2$, $|\phi_2|^2$ and $|\phi_3|^2$ all by the same amount. We need only note that for any solution, we must have at least two of the ϕ of order ξ in order to avoid a very large D -term of order $V_D \sim g^2 \xi^2$. Plugging these values in, we then find

$$\begin{aligned} V_D &\sim \mathcal{O}\left(\frac{\lambda^4 \xi^2}{g^2}\right) \\ V_F &\sim \mathcal{O}(\lambda^2 \xi^2). \end{aligned} \tag{5}$$

Note that the F -term generically does not vanish. There is one complex D -flat direction, but 6 real F -term equations which must be satisfied. So generally these terms cannot cancel. Furthermore, there is no runaway direction in the open-string configuration space which can restore supersymmetry. The only runaway direction for the D -term potential is $|\phi_1| \sim |\phi_2| \sim |\phi_3| \rightarrow \infty$. But in this limit it is clear that any linear or quadratic terms in the superpotential are dominated by the Yukawa term, and this leads to an F -term potential

$$V_F = \lambda^2 (|\phi_1|^2 |\phi_2|^2 + |\phi_2|^2 |\phi_3|^2 + |\phi_3|^2 |\phi_1|^2), \tag{6}$$

which blows up. Thus, in our expected limit $\lambda \ll g$, SUSY-breaking is dominated by the F -term. Moreover we have $F_\phi \sim \lambda \xi$, so even for a generic choice of ξ , the SUSY-breaking scale will be exponentially small.

The upshot of this section does not rely on the precise form of the superpotential, merely on the fact that there are fewer D -flat directions than F -term equations, and that all superpotential terms are non-perturbatively small.

Note that this mechanism is inherently stringy. From the point of view of effective field theory, there is no reason for λ to be small, and if instead we had $\lambda \sim \mathcal{O}(1)$ then we would have $F \sim \xi$. It is the structure of intersecting brane models which determines that the Yukawa couplings arise from suppressed worldsheet instantons, thus generating an exponentially small supersymmetry-breaking scale.

3.1 A vectorlike example

One can find a similar model with vectorlike matter. Suppose that some high-scale dynamics causes ϕ_3 to get a vev, breaking the gauge group $U(1)_b - U(1)_c$. In this case, we would be left at low energies with two gauge groups, $U(1)_a$ and $U(1)_b + U(1)_c$, the sum of which decouples. The matter content would be vectorlike and we would be left with essentially the Fayet model of [6].

$$\begin{aligned} V_D &= \frac{1}{2}g^2(|\phi_1|^2 - |\phi_2|^2 - \xi)^2 \\ W &= m\phi_1\phi_2. \end{aligned} \tag{7}$$

But here $m = \lambda\langle\phi_3\rangle$. The exponential suppression of λ ensures that our dynamically generated SUSY-breaking scale $F \sim \lambda\langle\phi_3\rangle\sqrt{\xi}$ will be small, as in [6].

4 Conclusions

The scenario presented here has many features in common with that presented in [6, 12]. Whereas their model generated an exponentially low mass in the superpotential from a D-instanton, the model presented here generates an exponentially small Yukawa coupling in the superpotential from a world-sheet instanton. This effectively replaces their scale $F \sim m\sqrt{\xi}$ with our scale $F \sim \lambda\xi$.

Our scenario does not involve placing branes at a singularity, however. In Type IIA/B, the two common ways of embedding a SM-like gauge theory in a string model are either by branes at singularities or by intersecting brane models. While the AKS setup seems naturally suited for generating DSB in the former class, our scenario is naturally suited to the latter.

Interestingly, this seems to be a very general scenario which should be quite common in intersecting brane models. Although we used only three D-branes in the hidden sector, it is clear that a more general hidden sector with more branes and more scalars would do equally well. The basic point is simply that vevs of the scalar fields are controlled by the FI-terms in V_D , with a small correction (which scales as $\frac{\lambda^2\xi}{g^2}$) due to the F -terms. At the minimum of the full scalar potential, we thus find that the FI-term contribution is fully cancelled by “uncorrected” scalar vevs, yielding only the “correction” which goes as $V_D \sim \frac{\lambda^4\xi^2}{g^2}$. Meanwhile, gauge invariance prevents the appearance of a tree-level mass term for non-vectorlike matter. Since there are fewer D -flat directions than F -term equations, one generically expects $V_F \neq 0$. The scale of V_F is then set by the cubic Yukawa couplings (and perhaps other instantons), which generate $V_F \sim \lambda^2\xi^2$ with much smaller higher-order corrections. As long as the size of the worldsheet instantons is large in string units (which is expected in the limit of weak coupling and large volume compactifications, which is best studied), a low dynamical scale should be generated naturally.

One way to think of the generality of this scenario is the following. If there are N gauge groups in the hidden sector with generic non-zero FI-terms, then one must give vevs to at least $\mathcal{O}(N)$ scalar fields (each charged under two gauge groups) in order to make the D -term potential small. In general, there will be at least two fields (oppositely) charged under each gauge group which get non-zero vevs. So for a gauge group $U(1)_G$, these scalars will be ρ_1 with charge +1 under $U(1)_G$ and -1 under some $U(1)_a$, plus another scalar ρ_2 with charge -1 under $U(1)_G$ and +1 under some $U(1)_b$. Now if $I_{ab} > 0$, then there will exist at least one scalar ρ_3 with charge +1 under $U(1)_a$ and charge -1 under $U(1)_b$. One can then write a Yukawa coupling $W = \lambda\rho_1\rho_2\rho_3$ which performs the task of breaking supersymmetry at a scale set by the exponentially small coupling λ (arising from a worldsheet instanton stretching between branes a , b and g). Or course, one can easily have $I_{ab} < 0$ (indeed, this feature was used in [13] to generate a flat inflationary potential). However, to ensure that there are no Yukawa terms which generate a non-zero F -term on the D -flat direction will require multiple fine-tunes on the signs of intersection numbers; the generic brane configuration will yield at least some non-zero F -term contribution, and this will be of order $\sim \mathcal{O}(\lambda^2\xi^2)$.

It is also interesting to note that in IBMs, supersymmetry breaking in the open-string sector is naturally mediated to visible sector by gauge-mediation. This is the case because hidden sector (including the SUSY-breaking sector) branes generically have non-trivial topological intersection with visible sector branes. This implies the existence of chiral multiplets charged under both SM and hidden sector gauge groups which can act as messengers. The dynamical generation of very small F -terms (perhaps TeV scale) in a hidden sector would

thus naturally fit into this scenario of gauge-mediation to the visible sector.

In constructions involving branes at singularities, it is more natural to have supersymmetry breaking mediated to the visible sector by gravity, although there are limits where gauge mediation is a better description[14]. In gravity-mediated scenarios one must usually do some work to ensure that undesirable FCNC's are avoided[15], while gauge mediation naturally avoids this problem for IBMs. On the other hand, for example, gauge unification is likely more easily understood for branes at singularities than for intersecting brane models. In some sense, these two methods of realizing the Standard Model display one characteristic reminiscent of dualities, namely, that nice features which are easy to understand in one model are difficult to understand in the other. As more avenues for dynamical supersymmetry-breaking are discovered for both intersecting brane models and branes at singularities, it will be interesting to discover how the two classes of models are related.

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