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Exponentially Small Supersymmetry Breaking from Extra Dimensions

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The supersymmetric "shining" of free massive chiral superfields in extra dimensions from a distant source brane can trigger exponentially small supersymmetry breaking on our brane of order $e^{-2\pi R}$, where R is the radius of the extra dimensions. This supersymmetry breaking can be transmitted to the superpartners in a number of ways, for instance by gravity or via the standard model gauge interactions. The radius R can easily be stabilized at a size $O(10)$ larger that the fundamental scale. The models are extremely simple, relying only on free, classical bulk dynamics to solve the hierarchy problem.

1 The four forces of nature are each characterized by a mass scale: $\sqrt{1/G_N} = M_P \approx 10^{19}$ GeV for gravity, $\Lambda_W \approx 10^3$ GeV for the weak interaction, $\Lambda_{QCD} \approx 0.1$ GeV for the strong interaction and $m_\gamma = 0$ for the electromagnetic interaction. What is the origin of these diverse scales? Over the last 25 years a single dominant viewpoint has developed: the largest scale, that of gravity, is fundamental, and the other scales are generated by a quantum effect in gauge theories known as dimensional transmutation. If the coupling strengths of the other forces have values $\alpha_P \approx 1/30$ at the fundamental scale, then a logarithmic evolution of these coupling strengths with energy leads, in non-Abelian theories, to the generation of a new mass scale

$$\Lambda \approx M_P\, e^{-1/\alpha_P} \tag{1}$$

where the interaction becomes non-perturbative. On the other hand, Abelian theories, like QED, remain perturbative to arbitrarily low scales. For strong and electromagnetic interactions this viewpoint is immediately successful; but for the weak interaction the success is less clear, since the weak interactions are highly perturbative at the scale Λ_W. If Λ_W is generated by a dimensional transmutation, it must happen indirectly by some new force getting strong and triggering the breakdown of electroweak symmetry. There have been different ideas about how this might occur: the simplest idea is technicolor, a scaled up version of the strong force [1]; another possibility has the new strong force first triggering supersymmetry breaking which in turn triggers electroweak symmetry breaking [2]. For our purposes the crucial thing about these very different schemes is that they have a common mechanism underlying the origin of Λ_W: a dimensional transmutation, caused by the logarithmic energy evolution of a gauge coupling constant, generates the exponential hierarchy of (1).

In this letter, we propose an alternative mechanism for generating Λ_W exponentially smaller than the fundamental scale. Our scheme requires two essential ingredients beyond the standard model: supersymmetry, and compact extra dimensions of space. The known gauge interactions reside on a 3-brane, and physics of the surrounding bulk plays a crucial role in generating an exponentially small scale of supersymmetry breaking.

Our mechanism is based on the idea of "shining" [3]. A bulk scalar field, ϕ, of mass m, is coupled to a classical source, J, on a brane at location $y = 0$ in the bulk, thereby acquiring an exponential profile $\phi \propto Je^{-m|y|}$ in all regions of the bulk distant from the source, $m|y| \gg 1$. If our brane is distant from the source, then this small exponential, arising from the propagation of the heavy scalar across the bulk, can provide an origin for very small dimensionless numbers on our brane, in particular for supersymmetry and electroweak symmetry breaking

$$\Lambda_W \propto M_*\, e^{-mR} \tag{2}$$

where R is the distance scale of our brane from the source brane, and M_* is the fundamental scale of the theory. The possibility of such a supersymmetry-breaking mechanism has been noted before qualitatively [3]. If some of the extra dimensions are very large, M_* can be significantly below M_P, and could even be of order Λ_W, providing an alternative viewpoint on the mass scales of the four fources of nature [4]. We are concerned with the case of $M_* \gg \Lambda_W$, although M_* need not be as large as M_P. In this letter we give an explicit construction of shining which preserves 4-dimensional supersymmetry, but triggers an exponentially small amount of supersymmetry breaking due to the presence of our brane. A possible worry is that R might run to infinity, thus minimizing the vacuum energy and restoring supersymmetry. We exhibit simple mechanisms, based on the same supersymmetric shining, which stabilize the extra dimensions with finite radius.

2 We begin by constructing a 5d theory, with a source brane shining an exponential profile for a bulk scalar, such that the equivalent 4d theory is exactly supersymmetric. The 5d theory possesses N=1 supersymmetry in a representation containing two scalar fields, ϕ and ϕ^c, together with a four-component spinor $\Psi = (\psi, \psi^c)$. The equivalent 4d theory has two families of chiral superfields $\Phi(y) = \phi(y) + \theta\psi(y) + \theta^2 F(y)$ and $\Phi^c(y) = \phi^c(y) + \theta\psi^c(y) + \theta^2 F^c(y)$. In the 4d theory, y can be viewed as a parameter labelling the families of chiral superfields.

Using this 4d chiral superfield notation, we write the bulk action as

$$S_B = \int d^4x\, dy \; \Big(\int d^4\theta (\Phi^\dagger \Phi + \Phi^{c\dagger}\Phi^c) + \int d^2\theta \Phi^c (m + \partial_y)\Phi \Big) \tag{3}$$

Viewed as a 4d theory, we have manifest supersymmetry, with the y integral summing over the family of chiral superfields. The form of the superpotential appears somewhat unusual; however, on eliminating the auxiliary fields, the action in terms of component fields describes a free Dirac fermion and two complex scalar fields in 5d. The 5d Lorentz invariance is not manifest in (3), but this form is useful to us, since it makes the 4d supersymmetry manifest.

Next we locate a 3-brane at $y = 0$, and require that it provides a source, J, for a chiral superfield in a way which preserves 4d supersymmetry:

$$W_S = \int dy \delta(y)\, J\Phi^c, \tag{4}$$

where we choose units so that the fundamental scale of the theory $M_* = 1$. The conditions that this source shines scalar fields into the bulk such that supersymmetry is not spontaneously broken are

$$F(y) = (m - \partial_y)\phi^c = 0 \tag{5}$$
$$F^c(y) = J\delta(y) + (m + \partial_y)\phi = 0 \tag{6}$$

The first of these does not have any non-trivial solutions that do not blow up at infinity, or which are well-defined on a circle. The second, however, has the solution

$$\phi(y) = -\theta(y) J e^{-my}, \tag{7}$$

in infinite flat space and

$$\phi(y) = \frac{-Je^{-my}}{1 - e^{-2\pi m R}} \quad y \in [0, 2\pi R), \tag{8}$$

on a circle. Thus we see that ϕ has taken on a non-zero profile in the bulk, but in a way that the energy of the system remains zero and one supersymmetry remains unbroken. Interestingly, this is not the profile that occurs with non-supersymmetric shining, but is asymmetric, shining in only one direction. One may have thought that the gradient energy for any profile of a bulk scalar field would neccessarily break supersymmetry, but our example shows this is not the case. The $|F^c|^2$ contribution to the vacuum energy includes the $|\partial_y\phi|^2 + |m\phi|^2$ terms as expected, but these are cancelled by $\phi^*\partial_y\phi$ terms, and at $y = 0$ by terms which arise because J is coupled to the combination $(m + \partial_y)\phi(0)$. Note that if we had written a linear term for Φ instead of Φ^c, we would have shined a profile for ϕ^c in the opposite direction. Likewise, if we had chosen a negative value for m, we would shine ϕ in the opposite direction, since the 5d theory is invariant under $m \to -m$, $y \to -y$.

3 Having learned how to shine a chiral superfield from a source brane across the bulk, we now investigate whether a probe brane, located far from the source at $y = \overline{y}$, can sample the small value of $\phi(\overline{y})$ to break supersymmetry by an exponentially small amount on the probe brane. In addition to superfields which contain the standard model fields, the probe brane contains a standard model singlet chiral superfield X, and has a superpotential

$$W_P = \int dy\, \delta(y - \overline{y})(W_{MSSM} + \Phi X) \tag{9}$$

where W_{MSSM} is the superpotential of the minimal supersymmetric standard model. This superpotential has F-flatness conditions

$$F^c(y) = J\delta(y) + (m + \partial_y)\phi = 0 \tag{10}$$
$$F(y) = \delta(y - \overline{y})x + (m - \partial_y)\phi^c \tag{11}$$
$$F_X = \phi(\overline{y}). \tag{12}$$

The first equation can only be satisfied by having a shined value for $\phi(\overline{y}) \neq 0$. Clearly, the first and third equations cannot be simultaneously satisfied: we have an O'Raifeartaigh theory, and supersymmetry is spontaneously broken. As always in an O'Raifeartaigh theory, at tree level there is a flat direction: the value for x is undetermined, and if it is non-zero it acts as a source shining ϕ^c. It is simple to understand what is going on. In the presence of the source brane, the field ϕ is shined from the source brane, generating an exponentially small linear term for X on the probe brane. After we have integrated out the heavy fields ϕ and ϕ^c we are simply left with the superpotential on the probe brane

$$W_P \sim Je^{-m\overline{y}}X, \tag{13}$$

which generates a nonzero $F_X \sim Je^{-m\overline{y}}$.

This is not a precise equality, as the probe brane resists a non-zero $\phi(\overline{y})$, and provides a back reaction on the bulk. It is simple to show that this effect is qualitatively insignificant.

If the fifth dimension is a circle, then we can imagine that the probe brane is stabilized at some location on the circle, or that it will drift such that it is immediately next to the source brane where the resulting supersymmetry breaking is smallest, as in figure 1. In either case, we generate an exponentially small supersymmetry breaking scale F_X.

Notice that this is *not* in the same spirit as recent works that use bulk dynamics to transmit distantly broken supersymmetry [5]. Rather, in our case, in the absence of either source or probe brane, supersymmetry remains unbroken. It is the simultaneous presence of

both branes that leads to the exponentially small supersymmetry breaking. A simple option for mediating the supersymmetry breaking from F_X to the standard model superpartners is to add non-renormalizable operators to the probe brane

$$\Delta S_P = \int d^4x dy \delta(y-\overline{y})\left(\int d^4\theta\left(\frac{1}{M_*^2}X^\dagger X Q^\dagger Q + ...\right)\right.$$
$$\left. + \int d^2\theta\left(\frac{1}{M_*}XW^\alpha W_\alpha + ...\right)\right) \quad (14)$$

where Q is a quark superfield and W^α a standard model gauge field strength superfield. We have inserted M_* explicitly, so that the soft masses of the standard model superpartners and x are $\widetilde{m} \sim F_X/M_* \sim (J/M_*)e^{-m\overline{y}}$. Until now we have not specified the values for J and m; the most natural values are $J \approx M_*^2$ and $m \approx M_*$.

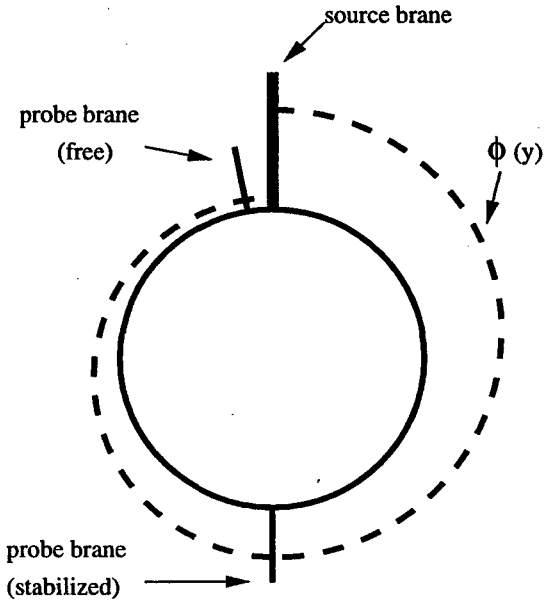

FIG. 1. The schematic profile of ϕ in the extra dimension. Whether our brane is stabilized at some position or free to move under the given forces, we can achieve an exponentially small value for ϕ and hence exponentially suppressed supersymmetry breaking.

Our entire theory is remarkably simple, and is specified by the bulk action S_B of (3), the source brane superpotential W_S of (4), and the interactions of (9) and (14) on our brane.

4 Mechanisms for dynamical supersymmetry breaking by dimensional transmutation [6] typically suffer from the "dilaton runaway problem" when embedded in string theory [7]: since the coupling constant α_P is a dynamical field, the vacuum energy is minimized as $\alpha_P \to 0$, where the theory becomes free. In our case, it appears there is an analogous problem. Taking the supersymmetry-breaking brane to be free to drift, the vacuum energy of the theory is

$$E \sim J^2 e^{-4\pi R m}, \quad (15)$$

so it is energetically favorable for the radius to grow to infinity. However, in contrast with dynamical supersymmetry breaking scenarios, where one must simply assume that the dilaton vev is somehow prevented from running to infinity, stabilizing R turns out to be quite simple.

Consider adding to the model of the previous section a second bulk multiplet (Φ', Φ'^c), of mass m', with interactions

$$W' = \int dy\ [\delta(y)J'\Phi'^c + \delta(y-\overline{y})X'(\Phi' + A)] \quad (16)$$

where A and J' are constants and X' is a chiral superfield. The terms in this superpotential are nearly identical to those of (9) and (4), except for the presence of the constant A on the probe brane. We assume that both A and J' are real. In complete analogy with the shining of ϕ, the scalar ϕ' acquires a profile

$$\phi'(y) = -J'\theta(y)e^{-m'y}. \quad (17)$$

Writing $\overline{y} = \theta R$, the F-flatness condition for X' becomes

$$m'R\theta = \log\frac{J'}{A(1-e^{-2\pi R m'})}, \quad (18)$$

which defines a real function $R(\theta)$ provided that $J'/A > 0$. We assume m' is less than m (by a factor of roughly 30, for very large M_*), so that, for a given value of θ, the radius is essentially determined by the condition $F_{X'} = 0$, with a small correction $\frac{\Delta R}{R} \sim \frac{m}{m'}e^{-m/m'}$ coming from the $|F_X|^2$ contribution to the potential. However, we have already seen that the vacuum energy is minimized when the probe brane drifts completely around the circle. The value of R is thus immediately fixed by equation (18), with $\theta = 2\pi$. Its precise value depends on A and J', but if we take their ratio to be of order unity, then we find $2\pi R m' \sim 1$. The supersymmetry breaking F-term is then $F_X \sim Je^{-2\pi m R} \sim Je^{-m/m'}$, so that the higher dimension interactions of (14) give superpartner masses

$$\widetilde{m} \sim e^{-m/m'}M_*. \quad (19)$$

In this model the mass of the radion, the field associated with fluctuations of the size of the circle, is $m_{radion} \sim F_X/M_P \sim 1$ TeV (M_*/M_P).

Alternatively one can stabilize R in an entirely supersymmetric fashion. Here we describe just one of a number of ways in which this can be done. Imagine supplementing the "clockwise" shining of ϕ' due to W' with "counterclockwise" shining of a different scalar $\widetilde{\phi}^c$ of comparable mass, $\widetilde{m}$, through the added superpotential terms

$$\widetilde{W} = \int dy\ \left[\delta(y)\widetilde{J}\widetilde{\Phi} + \delta(y-\overline{y})\widetilde{X}\left(\widetilde{\Phi}^c + B\right)\right]. \quad (20)$$

Note that because $\widetilde{\Phi}$ (rather than $\widetilde{\Phi}^c$) couples to the source, the shining is in the opposite direction as that of ϕ'. The F-flatness condition for $\widetilde{X}$,

$$\widetilde{m}R(2\pi-\theta)=\log\frac{B}{\widetilde{J}(1-e^{-2\pi R\widetilde{m}})}, \tag{21}$$

and the F-flatness condition for X' independently determine R as a function of θ, and for broad ranges of parameters the combined constraints are satisfied by unique values of θ and R. This supersymmetric stabilization of the radius yields $m_{radion} \sim M_*^2/M_P$, far above the TeV scale.

5 We have presented a complete model in which exponentially small supersymmetry breaking is generated as a bulk effect and communicated to the standard model via higher-dimension operators. It is straightforward to modify the model so that the supersymmetry breaking is mediated instead by gauge interactions [8].

Consider the O'Raifeartaigh superpotential

$$W = X(Y^2-\mu^2)+mZY. \tag{22}$$

At tree level x is a flat direction, but provided $\mu^2 < m^2/2$, radiative effects stabilize x at the origin and give $m_x^2 \sim \mu^2/16\pi^2$. Supersymmetry is broken by $F_X = -\mu^2$. Models using an O'Raifeartaigh superpotential to achieve low-energy supersymmetry breaking have been constructed in the past, but have required a small value for μ^2 to be input by hand. Instead, we use supersymmetric shining as an origin for the parameters μ^2 and m by coupling the brane superfields X, Y, and Z to the shone Φ according to

$$W_{hidden} = \lambda_1 X(Y^2-\Phi(\overline{y})^2)+\lambda_2\Phi(\overline{y})ZY, \tag{23}$$

where λ_1 and λ_2 are both of order unity and $\lambda_1 < \lambda_2^2/2$. Next we introduce couplings to messenger fields Q and $\overline{Q}$ transforming under the standard model gauge group [9],

$$W_{messenger} = \alpha_1 XQ\overline{Q}+\alpha_2\Phi(\overline{y})Q\overline{Q}. \tag{24}$$

By taking $\alpha_2^2 > \alpha_1\lambda_1$ we ensure that the messenger scalars do not acquire vevs. These superpotentials give Q and $\overline{Q}$ supersymmetric masses and supersymmetry-breaking mass splittings of comparable order, $M \sim \sqrt{F} \sim \phi(\overline{y})$. The messengers then feed the supersymmetry breaking into the standard model in the usual way, yielding soft supersymmetry-breaking parameters of order $\widetilde{m} \sim \frac{1}{16\pi^2}\phi(\overline{y})$. Fixing the radius R by either of the mechanisms already described then leads to $\widetilde{m} \sim \frac{M_*}{16\pi^2}e^{-m/m'}$. Note that this is truly a model of low-energy supersymmetry breaking, with $\sqrt{F} \sim 16\pi^2\widetilde{m} \sim 100$ TeV, allowing for decays of the NLSP within a detector length. Moreover, this small value for $\sqrt{F}$ is favored by cosmology in that it suppresses the gravitino energy density [10].

While there is typically a severe μ problem in gauge-mediated theories [11], it is easily solved with our mechanism by shining μ in the superpotential with a term

$$W \supset \lambda\phi(\overline{y})H_1H_2. \tag{25}$$

With $\lambda \sim 1/30$, problems of naturalness are much less severe than in theories where supersymmetry is broken dynamically. If $B\mu = 0$ at tree level, radiative effects can generate a small $B\mu$ and large $\tan\beta$ [12]. Likewise, in gravity mediated theories, a shined term $\int d^2\theta\Phi(\overline{y})H_1H_2$ can also generate an appropriate value for μ, while $\int d^4X^\dagger XH_1H_2$ generates $B\mu$. Although ϕ is related to supersymmetry breaking, this is distinct from the Giudice-Masiero mechanism. Absent the superfield X, supersymmetry is preserved, but the value of μ is unchanged.

Depending on whether supersymmetric or supersymmetry breaking stabilization of the radius is employed, the radion mass is either $m_{radion} \sim M_*^2/M_P$ or $m_{radion} \sim \sqrt{F}/M_P \sim 1$ eV (M_*/M_P). Even the latter case is safe, since the limit on the radion mass is on the mm^{-1} scale, at the limits of experimental probes of gravity at short distances.

6 Dimensional transmutation, (1), and shining, (2), are alternative mechanisms for taking a dimensionless input of order 30 and generating an exponentially small mass hierarchy. These mass hierarchies can explain the scales of symmetry breaking, for instance of a global flavor symmetry, or of supersymmetry, as we have discussed. While dimensional transmutation is a quantum effect requiring an initial coupling which is highly perturbative, $1/\alpha_P \approx 30$, shining is classical and requires a bulk distance scale of size $R \approx 30M_*^{-1}$. Such a radius can in turn be stabilized in a simple way. We presented two standard ways of communicating this exponentially small supersymmetry breaking, through higher-dimensional operators or via standard model gauge interactions. It is clearly possible to employ other mechanisms, such as those discussed in [5]. Our theories are remarkably simple, using only free classical dynamics in one extra dimension. Extensions to more dimensions should be straightforward. While we have concentrated on constructing effective theories with exponentially small global supersymmetry breaking, it will be interesting to embed these models in a consistent local supergravity. It will also be interesting to explore whether any of these mechanisms can be realized in the D-brane construction of non-BPS states in string theory.

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