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Supersymmetry without supersymmetry

Hock-Seng Goh, Markus A. Luty and Siew-Phang Ng

Department of Physics, University of Maryland
College Park, Maryland 20742, U.S.A.
E-mail: hsgoh@physics.umd.edu, luty@umd.edu, spng@physics.umd.edu

Abstract: We investigate the possibility that supersymmetry is not a fundamental symmetry of nature, but emerges as an accidental approximate global symmetry at low energies. This can occur if the visible sector is non-supersymmetric at high scales, but flows toward a strongly-coupled superconformal fixed point at low energies; or, alternatively, if the visible sector is localized near the infrared brane of a warped higher-dimensional spacetime with supersymmetry broken only on the UV brane. These two scenarios are related by the AdS/CFT correspondence. In order for supersymmetry to solve the hierarchy problem, the conformal symmetry must be broken below 10^{11} GeV. Accelerated unification can naturally explain the observed gauge coupling unification by physics below the conformal breaking scale. In this framework, there is no gravitino and no reason for the existence of gravitational moduli, thus eliminating the cosmological problems associated with these particles. No special dynamics is required to break supersymmetry; rather, supersymmetry is broken at observable energies because the fixed point is never reached. In 4D language, this can be due to irrelevant supersymmetry breaking operators with approximately equal dimensions. In 5D language, the size of the extra dimension is stabilized by massive bulk fields. No small input parameters are required to generate a large hierarchy. Supersymmetry can be broken in the visible sector either through direct mediation or by the F term of the modulus associated with the breaking of conformal invariance.

Keywords: Supersymmetry Breaking, Supersymmetry Phenomenology.

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1. Introduction

If supersymmetry (SUSY) solves the hierarchy problem, it implies the presence of a new spacetime symmetry in nature. How does this symmetry arise? The standard paradigm is that SUSY is an exact symmetry in the fundamental UV theory, and is broken spontaneously in the IR. In this paper, we consider the alternative possibility that UV physics is completely non-supersymmetric, and SUSY emerges as an accidental symmetry in the IR.[1]

An accidental symmetry arises when the UV theory has no relevant or marginal operators that can be added to the lagrangian to break the symmetry. In that case, all symmetry breaking effects flow to zero in the IR, and the theory becomes invariant under the symmetry at low energies even if the fundamental theory violates the symmetry maximally. A famous example is baryon and lepton number in the standard model. However, in weakly-coupled theories with scalars, scalar mass terms are always relevant, so SUSY cannot emerge as an accidental symmetry in such theories.

The situation can be very different in strongly-coupled theories. Suppose that there exists a strongly coupled superconformal theory without any relevant or marginal SUSY breaking operators that can be added to the lagrangian.[2] This fixed point will be attractive to all perturbations, so the boundary of the basin of attraction of the fixed point will consist of theories that have no approximate SUSY. For example, $\mathcal{N} = 1$ SU(N) SUSY QCD with

[1] The idea that 'fundamental' symmetries such as Lorentz invariance might arise as accidental symmetries of IR fixed points has been previously considered by H.B. Nielsen and collaborators [1].

[2] Every 4D CFT has a conserved stress energy operator $T_{\mu\nu}$. The traceless part of $T_{\mu\nu}$ has dimension 4, but the trace T has an anomalous dimension. T is a SUSY breaking operator that can be added to the lagrangian, but we assume that it has a positive anomalous dimension so that it is irrelevant.

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F flavors has a strongly-coupled fixed point near the middle of the conformal window ($F \simeq 2N$) [2]. The scalar masss operators have an uncalculable scaling dimension that is known to be larger than the canonical dimension [3]. It is possible that the anomalous dimensions are large enough that scalar masses are irrelevant in this theory, in which case this theory with the addition of large scalar mass terms flows to a superconformal fixed point in the IR.

Of course, SUSY must be broken at low energies to account for the absence of superpartners of the observed particles. The standard paradigm of exact SUSY in the UV requires special structure to break SUSY in the observable sector near the weak scale, e.g. dynamical SUSY breaking. While many models of dynamical SUSY breaking are known (see ref. [4] for a review), these are very special theories and SUSY breaking is generally not robust against perturbations. The present framework offers an alternative in which the small scale of SUSY breaking in the visible sector is simply explained by the fact that the superconformal fixed point is not reached at low energies. SUSY is therefore broken explicitly rather than spontaneously; there is no Goldstino. There are several possible mechanisms that can prevent the approach to the fixed point. One possibility is that there is a relevant SUSY breaking operator with a small coefficient. This does not give an explanation of the smallness of the SUSY breaking, although the small parameter may be natural if the relevant term breaks a symmetry. In this paper, we consider the alternative possibility that the approach to the fixed point is prevented by *irrelevant* operators. This is very natural in superconformal theories that have a moduli space of vacua in the SUSY limit. In such theories, irrelevant SUSY breaking effects will generate a potential on the moduli space, and can stabilize the moduli away from the origin. If the two lowest-dimension operators have dimensions that are somewhat close together, this can stabilize the scale modulus (dilaton) at a scale that is exponentially small compared to the fundamental scale. This mechanism is very generic, and can naturally generate a large hierarchy without small input parameters or fine tuning.

In this framework, the low-energy degrees of freedom of the visible sector are the remnants of the superconformal sector below the conformal breaking scale Λ_{IR}, which is therefore the compositeness scale for the standard model matter and gauge particles. Since SUSY breaking is communicated to the visible sector only through irrelevant operators, SUSY is an approximate (but not exact) symmetry of the visible sector, i.e.SUSY breaking in the visible sector is naturally below the scale Λ_{IR}. Direct mediation of SUSY breaking gives rise to scalar masses, gaugino masses, A terms, and μ and $B\mu$ terms, all of the same size. Understanding the absence of squark mixing requires additional structure, as in minimal supergravity. An alternative is that the CFT dynamics generates an F term for the dilaton, the modulus associated with the scale of conformal symmetry breaking. This naturally gives SUSY breaking below the scale Λ_{IR} provided that the CFT has a small parameter that breaks $\text{U}(1)_R$ symmetry. In this case, SUSY breaking in the visible sector is naturally dominated by anomaly-mediated SUSY breaking [5]. If the visible sector is the minimal supersymmetric standard model, the slepton mass-squared terms are negative. However, completely realistic non-minimal models can be constructed in this framework, using e.g.the ideas of refs. [6, 7].

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An important general consequence of this framework is that the gravitational sector is completely non-supersymmetric. In particular, there is no gravitino in the spectrum. This is similar to recent models in which the gravitino mass is far above the weak scale [8, 9] (see also ref. [10]). In fact, the present framework can be thought of as a limit of the model of ref. [8] with a high SUSY breaking scale. The absence of the gravitino eliminates the constraint on the inflationary reheat temperature that comes from the condition that gravitinos are not overproduced. Furthermore, since fundamental physics is non-supersymmetric, there there are no gravitational moduli, scalar fields with Planck-suppressed couplings that are generally present in string theory and higher-dimensional SUSY theories. This is a very good thing, because gravitational moduli cause severe cosmological difficulties that cannot be 'inflated away', and have been viewed as a major obstacle to realistic string model-building (see e.g. [11]). The present models do have a dilaton modulus in the 4D CFT description, but this modulus couples with more than gravitational strength, and does not give rise to cosmological difficulties.

Because gravity is not supersymmetric, gravity loops will generate SUSY breaking in the visible sector. However, these loops will be cut off at the scale $\Lambda_{\rm IR}$ where the conformal symmetry and SUSY are restored. Above the scale $\Lambda_{\rm IR}$, the gravity loops generate a perturbation corresponding to an irrelevant operator, which is therefore suppressed by the superconformal dynamics [12]. In order for the visible scalar masses to be naturally of order 100 GeV, the scale of conformal symmetry breaking must be below 10^{11} GeV.

Because the standard model matter and gauge fields are composite below 10^{16} GeV, gauge coupling unification cannot take place in the usual way. However, accelerated unification [13] can easily lower the unification scale below the compositeness scale, thus explaining the observed unification of the standard model gauge couplings.

All of the important features of this model arise as direct consequences of strong conformal dynamics with no relevant operators. It is now well understood that 5D gravity theories in anti de Sitter (AdS) space provide 'dual' descriptions of 4D strongly-coupled conformal field theories [14, 15].[3] We can therefore write explicit 5D models that realize the framework described above. The 5D models are of the Randall-Sundrum (RS) type [16], where the UV brane breaks SUSY, while the bulk and the IR brane are supersymmetric. The visible sector is localized on the IR brane. The 5D description is weakly coupled, making explicit calculations possible. In particular, we can easily understand SUSY breaking in the visible sector. We will construct an explicit 5D model as an existence proof, but we stress that the main features are generic to superconformal theories with only irrelevant SUSY breaking operators.

This paper is organized as follows. In section 2, we present an explicit 5D RS model that realizes the ideas outlined above. We consider radius stabilization and construct the low-energy 4D effective field theory, which we use to analyze SUSY breaking in the visible

[3]5D AdS theories are invariant under SO(4, 2), the 4D conformal symmetry, so it is rigorously true that any 5D AdS theory describes some 4D conformal field theory.

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sector. We interpret our results in the language of 4D conformal field theories and argue that the basic features are very general. In section 3, we briefly discuss phenomenology, and section 4 contains our conclusions.

2. 5D model

2.1 Definition of the model

Our model is based on the Randall-Sundrum (RS) model [16]. This is a 5D spacetime compactified on a $S^1/\mathbb{Z}_2$ orbifold, with metric

$$ds^2 = e^{-2\sigma(y)}\eta_{\mu\nu}dx^\mu dx^\nu + dy^2\,, \tag{2.1}$$

where y is a periodic variable with period 2ℓ, and $\sigma(y)$ is a periodic function defined by

$$\sigma(y) = k|y| \qquad \text{for} \quad -\ell < y \le +\ell\,. \tag{2.2}$$

Note that this implies

$$\begin{aligned}\sigma' &= \frac{d\sigma}{dy} = k\,\mathrm{sgn}(y) \qquad \text{for} \quad -\ell < y \le +\ell,\\ \sigma'' &= \frac{d^2\sigma}{dy^2} = 2k\left[\delta(y) - \delta(y-\ell) + \cdots\right].\end{aligned} \tag{2.3}$$

The physical region is $0 \le y \le \ell$. The boundary at $y = 0$ is the 'UV brane' (or 'Planck brane') where the zero mode of graviton is localized, and the boundary at $y = \ell$ is the 'IR brane.' We assume that the physics of this model is controlled by a single fundamental scale $\Lambda_{\rm UV} \sim M_{\rm P}$. This means that we take all couplings in the action to be of order 1 in units of $\Lambda_{\rm UV}$. The effect of the metric eq. (2.1) is that physical scales on the IR brane are 'warped down' to the scale

$$\Lambda_{\rm IR} = \Lambda_{\rm UV}\omega\,, \tag{2.4}$$

where

$$\omega = e^{-k\ell} \tag{2.5}$$

is the 'warp factor.' There can be an exponentially large hierarchy between the fundamental scale and $\Lambda_{\rm IR}$, provided that the size of the extra dimension ℓ can be stabilized at a value somewhat larger than $k^{-1} \sim \Lambda_{\rm UV}^{-1}$, so that $\omega \ll 1$. This is the hierarchy generating mechanism of Randall and Sundrum [16].

The RS model is interesting in its own right, but an additional motivation to consider this model is that it can be interpreted as a strongly-coupled 4D conformal field theory (CFT) [15, 17, 18]. The origin of this equivalence is that the bulk AdS_5 geometry has a $\mathrm{SO}(4,2)$ symmetry, which is isomorphic to the 4D conformal group, which acts on the branes as a 4D conformal transformation. In this equivalence, the bulk Kaluza-Klein (KK) modes are identified with excitations of the CFT. The UV brane acts as a UV cutoff on these modes, while the IR brane gives rise to spontaneous breaking of the conformal

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invariance.[4] Bulk fields are associated with operators of the CFT. Scalar operators that are irrelevant (respectively relevant) are associated with scalar modes with bulk mass $m^2 > 0$ (respectively $m^2 < 0$).

We are interested in 4D CFT's where the UV physics breaks SUSY, but the theory flows toward a superconformal fixed point in the IR. This means we want a 5D model where the couplings on the UV brane break SUSY maximally, while the action for the bulk and the IR brane is supersymmetric. This is radiatively stable by 5D locality. We also want the CFT to have only irrelevant perturbations. This means that all scalar fields must have positive bulk masses. We will therefore add massive hypermultiplets in the bulk, which will play an important role in stabilization and SUSY breaking.[5]

At energies below the mass of the lightest KK mode $m_{\rm KK} \sim \Lambda_{\rm IR}$ the physics can be described by a 4D effective theory that is approximately supersymmetric. The light degrees of freedom consist of the $\mathcal{N} = 1$ SUGRA multiplet, the radion chiral multiplet

$$\omega = e^{-k\ell} + \cdots + \theta^2 F_\omega \,, \tag{2.6}$$

and the light fields localized on the IR brane. The effective lagrangian is [19]

$$\begin{aligned} \mathcal{L}_{4,\rm eff} = &-\frac{M_5^3}{k} \int d^4\theta\, (\omega^\dagger \omega - \varphi^\dagger \varphi) + \\ &+ \int d^4\theta\, \omega^\dagger \omega K_{\rm IR} + \left(\int d^2\theta\, \omega^3 W_{\rm IR} + \text{h.c.} \right) + \\ &+ \text{SUSY breaking terms}\,, \end{aligned} \tag{2.7}$$

where $\varphi = 1 + \theta^2 F_\varphi$ is the conformal compensator and the superspace integrals are shorthand for the the covariant F and D projections of the superconformal tensor calculus [20]. The 4D Planck scale is

$$M_{\rm P}^2 = \frac{M_5^3}{k}(1 - \omega^2) \simeq \frac{M_5^3}{k}. \tag{2.8}$$

$K_{\rm IR}$ and $W_{\rm IR}$ are the Kählerpotential and superpotential of the fields localized on the IR brane. Note that the field ω has a canonical kinetic term, and that it couples to physics on the IR brane as a dilaton. This is the nonlinear realization of the conformal symmetry, which will play an important role in what follows.

2.2 SUSY breaking from 5D SUGRA

We now begin our discussion of SUSY breaking on the IR brane (the visible sector). Any SUSY breaking effects must arise by communication with the UV brane via bulk modes. In this subsection, we discuss the effects of the 5D SUGRA fields. We will discuss the effects of the bulk hypermultiplets after we have discussed stabilization.

Tree-level SUGRA KK exchange does not give rise to SUSY breaking operators on the IR brane [19, 21, 22]. There is a potential tree-level SUSY breaking effect from a constant

[4] More precisely, the conformal symmetry is nonlinearly realized by the position of the IR brane in the limit where the UV brane is at infinity [18].

[5] The trace of the CFT stress-energy tensor corresponds to the 5D dilaton state. We therefore assume that the 5D dilaton is more massive than the hypermultiplets.

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superpotential on the IR brane. This generates a nonzero VEV for F_ω, which gives rise to anomaly-mediated SUSY breaking on the IR brane [8]. If the constant superpotential is order 1 in units of $\Lambda_{\rm UV}$, we obtain $F_\omega/\omega \sim \Lambda_{\rm IR}$, where the left-hand sideis the order parameter for anomaly mediation on the IR brane. In order to obtain SUSY breaking masses at the weak scale, we must have $\Lambda_{\rm IR} \lesssim 10\,{\rm TeV}$. Since $\Lambda_{\rm IR}$ is also the compositeness scale for the standard model gauge fields, this implies that the standard model gauge fields are strongly coupled below 10 TeV. This requires a large number of charged states below 10 TeV, and the masses of these extra states must be finely tuned to get the observed low-energy gauge couplings. To avoid this unattractive scenario, we assume that the constant superpotential term is absent or small, which is natural by $U(1)_R$ symmetry.

We now consider SUGRA loop contributions to SUSY breaking on the IR brane, e.g.scalar masses. For loop momenta below $m_{\rm KK} \sim \Lambda_{\rm IR}$, the loop diagram is identical to a 4D loop diagram with a graviton line. This integral is effectively cut off for momenta above $m_{\rm KK} \sim \Lambda_{\rm IR}$ by higher-dimensional locality. This gives [8]

$$\Delta m^2_{\rm scalar} \sim \frac{1}{16\pi^2}\frac{\Lambda^4_{\rm IR}}{M^2_{\rm P}}\,. \tag{2.9}$$

Demanding that this contribution to the scalar masses be of order 100 GeV or less gives

$$\Lambda_{\rm IR} \lesssim 10^{11}\ {\rm GeV}\,. \tag{2.10}$$

If $\Lambda_{\rm IR} \sim 10^{11}$ GeV, this gives a flavor universal contribution to the scalar masses of order 100 GeV. In models where SUSY breaking in the visible sector is anomaly-mediated, an additional positive scalar mass-squared contribution can make the slepton masses positive in the minimal supersymmetric standard model. It would therefore be extremely interesting to compute the sign of the scalar mass contribution eq. (2.9).[6] Gaugino masses and A terms from SUGRA loops are negligibly small.

2.3 Casimir energy

Because SUSY is broken, there will be a nonzero Casimir energy from states in the bulk that feel SUSY breaking via couplings to the UV brane. The Casimir energy depends on the radius, and therefore contributes to the radius potential.

The Casimir energy can be written as a sum over KK modes. (In the 4D CFT interpretation, these are bound states of the strongly-coupled CFT.) The radius-dependent part of the Casimir energy is UV finite, and is therefore dominated by the contribution from the lowest lying KK states, which are approximately supersymmetric. The SUSY violating mass splittings are due to gravitational strength interactions, and are therefore of order

$$\Delta m_{\rm KK} \sim \frac{m^3_{\rm KK}}{M^2_{\rm P}} \sim \omega^3\,. \tag{2.11}$$

The Casimir energy vanishes when the spectrum is supersymmetric, so we have

$$V_{\rm Casimir} \sim m^3_{\rm KK}\Delta m_{\rm KK} \sim \omega^6\,. \tag{2.12}$$

[6]For the calculation in flat 5D theories, see refs. [23, 24].

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This agrees with explicit calculations (see e.g.ref. [25]).[7] We will consider stabilization mechanisms such that this contribution to the potential is negligible, so we do not need to know the sign of the Casimir energy.

2.4 Radius stabilization

We now give a detailed discussion of radius stabilization in the 5D model outlined above. The 4D CFT interpretation of the radion is the scale modulus (dilaton) associated with spontaneous conformal symmetry breaking [18], so this is equivalent to dilaton stabilization in the CFT.

In the 5D description, the radion potential is generated by the Goldberger-Wise mechanism [26] with bulk scalar fields with positive mass-squared.[8] In the CFT description, the bulk scalars parameterize the effects of irrelevant operators in the CFT. For a single scalar, we expect a potential of the form

$$V_{\text{eff}} = a\omega^n \,, \tag{2.13}$$

with $n > 0$. By itself this will give a runaway potential, but we can obtain a stable minimum if there are several such terms:

$$V_{\text{eff}} = a_1\omega^{n_1} + a_2\omega^{n_2} \,, \tag{2.14}$$

with $n_1 > n_2 > 0$. We find a local minimum at a small value of ω if n_1 and n_2 are close in value. If we define $\epsilon = n_1 - n_2$, then

$$\omega = \left(-\frac{(n_1 - \epsilon)a_2}{n_1 a_1} \right)^{1/\epsilon} \tag{2.15}$$

is a minimum provided that $a_1 > 0$, $a_2 < 0$. (In fact, $V_{\text{eff}} < 0$ at the minimum, so this vacuum has lower than the asymptotic vacuum $\omega = 0$.) Note that ω is naturally exponentially small if the factor in parentheses is positive and less than one, and ϵ is moderately small. The potential eq. (2.14) has the same form as the modulus potential in 'racetrack' models [27]. The radion field ω has a kinetic term with coefficient M_{P}^2 (see eqs. (2.7) and (2.8)), so the physical mass of the radion is of order

$$m_{\text{radion}}^2 \sim \frac{\epsilon a_1}{M_{\text{P}}^2}\omega^{n_1-2} \,. \tag{2.16}$$

We now discuss in detail how potentials of this form can arise from bulk hypermultiplets. The action for a hypermultiplet in the RS background was given in terms of $\mathcal{N} = 1$ superfields by Martí and Pomarol in ref. [28]. We add a general SUSY breaking potential on the UV brane, and a superpotential on the IR brane. These are the leading brane-localized terms in a low-energy expansion. The action is therefore

$$S_{\text{hyper}} = \int d^4x \int_{-\ell}^{\ell} dy\, \mathcal{L}_5 \,, \tag{2.17}$$

[7] We thank A. Pomarol for helpful discussions on Casimir energy.

[8] ref. [26] considered scalars with $m^2 \ll k^2$, corresponding to almost marginal 4D CFT operators; we consider $m^2 \gtrsim k^2$.

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where[9]

$$\mathcal{L}_5 = \int d^4\theta\, e^{-2\sigma}(\Phi^\dagger\Phi + \tilde{\Phi}^\dagger\tilde{\Phi}) + \left[\int d^2\theta\, e^{-3\sigma}\left(\frac{1}{2}\tilde{\Phi}\overleftrightarrow{\partial}_y\Phi + c\sigma'\tilde{\Phi}\Phi\right) + \text{h.c.}\right] - \\ - \delta(y)U(\Phi, F) + \delta(y-\ell)\,\omega^3\left[\int d^4\theta\, W(\Phi) + \text{h.c.}\right]. \tag{2.18}$$

The AdS/CFT correspondence relates the mass of states in the bulk to the dimension of an operator in the CFT. A 5D scalar with bulk mass m corresponds to an operator of dimension $d = 2 + \sqrt{4 + m^2/k^2}$. The bulk mass of the scalars from the hypermultiplet action above is

$$m^2_{\Phi,\tilde{\Phi}} = k^2\left(c \mp \frac{3}{2}\right)\left(c \pm \frac{5}{2}\right). \tag{2.19}$$

so the dimensions of the operators associated with the scalar components of Φ and $\tilde{\Phi}$ are

$$\dim(\mathcal{O}_{\Phi,\tilde{\Phi}}) = 2 + \left|c \pm \frac{1}{2}\right|. \tag{2.20}$$

If we want the operators associated with both Φ and $\tilde{\Phi}$ to be irrelevant, we must have $|c| > 5/2$.

For a scalar ϕ of mass m corresponding to an operator of dimension d, the general solution to the bulk equations of motion is

$$\phi = Ae^{d\sigma} + Be^{(4-d)\sigma}. \tag{2.21}$$

The coefficients A and B are fixed by the boundary conditions. The CFT interpretation of the coefficients is as follows. A is associated with a VEV of the operator

$$A = \frac{\langle\mathcal{O}\rangle}{2d-4}, \tag{2.22}$$

while B is associated with adding a term to the CFT lagrangian

$$\Delta\mathcal{L}_{\text{CFT}} = \lambda\mathcal{O}, \tag{2.23}$$

with

$$B \propto \lambda. \tag{2.24}$$

For irrelevant operators ($d > 4$) the second term in eq. (2.21) is exponentially decreasing in the IR, and therefore the value of the coefficient B will be determined by physics on the UV brane. In models where all dimensionful couplings are order 1 in units of Λ_{UV}, we therefore expect $B \sim \Lambda_{\text{UV}}^{3/2}$. The first term in eq. (2.21) grows in the IR, and is therefore determined by physics on the IR brane. The scale of physics on the IR brane is set by Λ_{IR}, so this will generally fix $A \sim \Lambda_{\text{IR}}^{3/2} \ll \Lambda_{\text{UV}}^{3/2}$. We therefore say that the first (second) term in eq. (2.21) is IR (UV) dominated, respectively.

[9] We write the action in terms of the two-sided derivative $\tilde{\Phi}\overleftrightarrow{\partial}_y\Phi = \tilde{\Phi}\partial_y\Phi - (\partial_y\tilde{\Phi})\Phi$. This is without loss of generality, since $f\tilde{\Phi}\overleftrightarrow{\partial}_y\Phi = 2f\tilde{\Phi}\partial_y\Phi + \partial_y f\tilde{\Phi}\Phi + \text{total derivative}$

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We now solve the equations of motion. We look for solutions depending only on y. The $\tilde{\Phi}$ equation of motion is

$$\partial_y F + \left(c - \frac{3}{2}\right)\sigma' F = 0\,. \tag{2.25}$$

Imposing the condition that F is periodic and even, the most general solution is

$$F = F_0 e^{-(c-3/2)\sigma}\,, \tag{2.26}$$

where F_0 is a constant of integration. We can define

$$\begin{aligned} F_{\rm UV} &= F(0) = F_0\,, \\ F_{\rm IR} &= F(\ell) = F_0 \omega^{c-3/2}\,. \end{aligned} \tag{2.27}$$

The Φ equation of motion is

$$e^{-3\sigma}\partial_y \tilde{F} - \left(c + \frac{3}{2}\right)\sigma' e^{-3\sigma}\tilde{F} = -\delta(y)\frac{\partial U}{\partial \Phi} + \delta(y-\ell)\omega^3 \frac{\partial^2 W}{\partial \Phi^2} F\,. \tag{2.28}$$

Imposing the condition that F is periodic and odd, the most general solution is

$$\tilde{F} = \tilde{F}_0 \frac{\sigma'}{k} e^{(c+3/2)\sigma}\,, \tag{2.29}$$

where $\tilde{F}_0$ is a constant of integration. Even though $\tilde{F}$ is odd under the orbifold Z_2, it is discontinuous at the boundaries, so it effectively has a nonvanishing value on each boundary. It is convenient to define

$$\begin{aligned} \tilde{F}_{\rm UV} &= \lim_{y\to 0+} \tilde{F} = \tilde{F}_0\,, \\ \tilde{F}_{\rm IR} &= \lim_{y\to \ell-} \tilde{F} = \tilde{F}_0 \omega^{-(c+3/2)}\,. \end{aligned} \tag{2.30}$$

In terms of these quantities, the jump conditions at the UV and IR branes are

$$\tilde{F}_{\rm UV} = -\tfrac{1}{2}\frac{\partial U}{\partial \Phi_{\rm UV}}\,, \tag{2.31}$$

$$\tilde{F}_{\rm IR} = -\tfrac{1}{2}\frac{\partial^2 W}{\partial \Phi_{\rm IR}^2} F_{\rm IR}\,, \tag{2.32}$$

where $\Phi_{\rm UV} = \Phi(0)$, $\Phi_{\rm IR} = \Phi(\ell)$.

The $\tilde{F}$ equation of motion is

$$e^{-3\sigma}\partial_y \Phi + \left(c - \frac{3}{2}\right)\sigma' e^{-3\sigma}\Phi + e^{-2\sigma}\tilde{F}^\dagger = 0\,. \tag{2.33}$$

Imposing the condition that Φ is periodic and even, the most general solution is

$$\Phi = \Phi_0 e^{-(c-3/2)\sigma} - \frac{\tilde{F}_0^\dagger}{(2c+1)k} e^{(c+5/2)\sigma}\,, \tag{2.34}$$

where Φ_0 is a constant of integration.

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Finally, the F equation of motion is

$$e^{-3\sigma}\partial_y\tilde{\Phi} - \left(c+\frac{3}{2}\right)\sigma' e^{-3\sigma}\tilde{\Phi} - e^{-2\sigma}F^\dagger = -\delta(y)\frac{\partial U}{\partial F} + \delta(y-\ell)\omega^3\frac{\partial W}{\partial \Phi}. \tag{2.35}$$

Imposing the condition that $\tilde{\Phi}$ is periodic and odd, the most general solution is

$$\tilde{\Phi} = \frac{\sigma'}{k}\left[\tilde{\Phi}_0 e^{(c+3/2)\sigma} - \frac{F_0^\dagger}{(2c-1)k}e^{-(c-5/2)\sigma}\right], \tag{2.36}$$

where $\tilde{\Phi}_0$ is a constant of integration. To write the boundary conditions, we define

$$\begin{aligned}\tilde{\Phi}_{\rm UV} &= \lim_{y\to 0+}\tilde{\Phi} = \tilde{\Phi}_0 - \frac{F_0^\dagger}{(2c-1)k}\,,\\ \tilde{\Phi}_{\rm IR} &= \lim_{y\to \ell-}\tilde{\Phi} = \tilde{\Phi}_0\omega^{-(c+3/2)} - \frac{F_0^\dagger}{(2c-1)k}\omega^{c-\frac{5}{2}}\,.\end{aligned} \tag{2.37}$$

The jump conditions at the UV and IR branes can then be written

$$\tilde{\Phi}_{\rm UV} = -\frac{1}{2}\frac{\partial U}{\partial F_{\rm UV}}\,, \tag{2.38}$$

$$\tilde{\Phi}_{\rm IR} = -\frac{1}{2}\frac{\partial W}{\partial \Phi_{\rm IR}}\,. \tag{2.39}$$

The fields $\tilde{\Phi}$ and $\tilde{F}$ are discontinuous at the boundaries. In our formulation with auxiliary fields, this is because the equations are first-order with delta function terms on the boundaries. Integrating out the auxiliary fields give rise to terms proportional to powers of delta functions, which naïvely are too singular to have a good continuum limit. However, supersymmetry and the orbifold projection evidently make sense out of these singular brane terms and give rise to the discontinuities at the boundaries.

To summarize, the solution is given by eqs. (2.26), (2.29), (2.34), (2.36). The four constants of integration F_0, $\tilde{F}_0$, Φ_0, and $\tilde{\Phi}_0$ are to be determined from the four discontinuity conditions in eqs. (2.31), (2.32), (2.38), and (2.39). Note that the jump conditions contain no explicit dependence on ω when expressed entirely in terms of UV or IR quantities.

We now use these results to write the effective potential. At the classical level, this is obtained by substituting the solutions to the equations of motion given above into the action and performing the integral over y to obtain a potential that depends on ω. Because the bulk terms are quadratic in Φ and $\tilde{\Phi}$, imposing the bulk equations reduces the effective potential to boundary terms. Using the jump conditions to simplify the result, we obtain the rather elegant expression

$$\begin{aligned}V_{\rm eff} &= U(\Phi_{\rm UV}, F_{\rm UV}) + \left(\Phi_{\rm UV}\tilde{F}_{\rm UV} + \tilde{\Phi}_{\rm UV}F_{\rm UV} + \text{h.c.}\right) +\\ &\quad + \omega^3\left(-\Phi_{\rm IR}\tilde{F}_{\rm IR} + \tilde{\Phi}_{\rm IR}F_{\rm IR} + \text{h.c.}\right).\end{aligned} \tag{2.40}$$

There is implicit dependence on ω through the boundary values of the fields.

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We now specialize to the case where $c > 5/2$, so that the scalar components of Φ and $\tilde{\Phi}$ are associated with operators of dimension

$$d = \dim(\mathcal{O}_\Phi) = c + \frac{5}{2}, \qquad \tilde{d} = \dim(\mathcal{O}_{\tilde{\Phi}}) = c + \frac{3}{2}, \tag{2.41}$$

with $d, \tilde{d} > 4$. We also assume that all couplings in the lagrangian are order one in units of $\Lambda_{\rm UV}$. Then we have

$$F = F_{\rm UV} e^{(4-d)\sigma}. \tag{2.42}$$

This is UV dominated, so we parameterize it by $F_{\rm UV}$. We also have

$$\tilde{F} = \frac{\sigma'}{k} \tilde{F}_{\rm IR} \omega^{\tilde{d}} e^{\tilde{d}\sigma}. \tag{2.43}$$

This is IR dominated, so we parameterize it by $\tilde{F}_{\rm IR}$. The solution for Φ is

$$\Phi = \left[\Phi_{\rm UV} + \frac{\tilde{F}^\dagger_{\rm IR}\omega^{\tilde{d}}}{(2d-4)k}\right] e^{(4-d)\sigma} - \frac{\tilde{F}^\dagger_{\rm IR}\omega^{\tilde{d}}}{(2d-4)k} e^{d\sigma}. \tag{2.44}$$

We will see that $\tilde{F}_{\rm IR} \ll \mathcal{O}(\omega)$ as a result of the jump equations, so the last term is small for all values of y. Φ is therefore UV dominated, and we parameterize it using $\Phi_{\rm UV}$. Finally, we have

$$\tilde{\Phi} = \frac{\sigma'}{k}\left[\left(\tilde{\Phi}_{\rm IR}\omega^{\tilde{d}} + \frac{F^\dagger_{\rm UV}}{(2\tilde{d}-4)k}\omega^{2\tilde{d}-4}\right) e^{\tilde{d}\sigma} - \frac{F^\dagger_{\rm UV}}{(2\tilde{d}-4)k} e^{(4-\tilde{d})\sigma}\right]. \tag{2.45}$$

$\tilde{\Phi}$ is unsuppressed at both the UV and IR branes. We parameterize it by $\tilde{\Phi}_{\rm IR}$.

The constants of integration $F_{\rm UV}$, $\tilde{F}_{\rm IR}$, $\Phi_{\rm UV}$, and $\Phi_{\rm IR}$ are determined by the jump equations. eqs. (2.31) and (2.32) are

$$\tilde{F}_{\rm IR} = -\frac{1}{2} F_{\rm UV}\omega^{d-4} \frac{\partial^2 W}{\partial \Phi^2_{\rm IR}}, \tag{2.46}$$

$$\frac{\partial U}{\partial \Phi_{\rm UV}} = -2\tilde{F}_{\rm IR}\omega^{\tilde{d}}. \tag{2.47}$$

eq. (2.46) implies that $\tilde{F}_{\rm IR} \lesssim \mathcal{O}(\omega^{d-4}) \ll \mathcal{O}(\omega)$, implying that Φ is UV dominated (see eq. (2.44)). The jump conditions eqs. (2.38) and (2.39) can be written

$$F^\dagger_{\rm UV} = \frac{(2\tilde{d}-4)k}{1-\omega^{2\tilde{d}-4}}\left[\frac{1}{2}\frac{\partial U}{\partial F_{\rm UV}} + \tilde{\Phi}_{\rm IR}\omega^{\tilde{d}}\right], \tag{2.48}$$

$$\tilde{\Phi}_{\rm IR} = -\frac{1}{2}\frac{\partial W}{\partial \Phi_{\rm IR}}. \tag{2.49}$$

From eqs. (2.46) through (2.49) we can see that there are generically solutions with $\Phi_{\rm UV}, F_{\rm UV} = \mathcal{O}(\omega^0)$. The leading approximation for $\Phi_{\rm UV}$ and $F_{\rm UV}$ is obtained by solving

$$0 = \frac{\partial U}{\partial \Phi_{\rm UV}}, \tag{2.50}$$

$$F^\dagger_{\rm UV} = (\tilde{d}-2)k\frac{\partial U}{\partial F_{\rm UV}}. \tag{2.51}$$

The corrections depend on the form of the IR superpotential.

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Let us consider an example where

$$W = \kappa \Phi \,. \tag{2.52}$$

We then have

$$\tilde{\Phi}_{\rm IR} = -\frac{1}{2}\kappa \,, \qquad \tilde{F} \equiv 0 \,, \tag{2.53}$$

while $\Phi_{\rm UV}$ and $F_{\rm UV}$ are determined by solving the equations

$$0 = \frac{\partial U}{\partial \Phi_{\rm UV}} \,, \tag{2.54}$$

$$F_{\rm UV}^\dagger = \frac{(\tilde{d}-2)k}{1-\omega^{2\tilde{d}-4}} \left(\frac{\partial U}{\partial F_{\rm UV}} - \kappa \omega^{\tilde{d}} \right) . \tag{2.55}$$

Expanding in powers of ω,

$$\Phi_{\rm UV} = \Phi_{\rm UV}^{(0)} + \omega^{\tilde{d}} \Phi_{\rm UV}^{(1)} + \mathcal{O}\left(\omega^{2\tilde{d}-4}\right) , \tag{2.56}$$

$$F_{\rm UV} = F_{\rm UV}^{(0)} + \omega^{\tilde{d}} F_{\rm UV}^{(1)} + \mathcal{O}\left(\omega^{2\tilde{d}-4}\right) , \tag{2.57}$$

we obtain the leading ω-dependent contribution to the effective potential:

$$V_{\rm eff} = \left[-\kappa F_{\rm UV}^{(0)} + \text{h.c.} \right] \omega^{\tilde{d}} + \mathcal{O}\left(\omega^{2\tilde{d}-4}\right) . \tag{2.58}$$

It is clear that the coefficient of $\omega^{\tilde{d}}$ can have either sign, as required for stabilization. If we want the Casimir energy to be negligible compared to the effects discussed here, we must have $\tilde{d} < 6$.

The results above are in agreement with the expectations of the AdS/CFT correspondence. The CFT operator of lowest dimension associated with the hypermultiplet is $\mathcal{O}_{\tilde{\Phi}}$, with dimension $\tilde{d}$. If we add to the CFT lagrangian a term

$$\Delta \mathcal{L}_{\rm UV} = \tilde{\lambda} \mathcal{O}_{\tilde{\Phi}} \,, \tag{2.59}$$

conformal invariance implies that the effective potential has the form

$$V_{\rm eff} = \Lambda_{\rm IR}^4 f(\tilde{\lambda}_{\rm eff}(\Lambda_{\rm IR})) \,, \tag{2.60}$$

where $\tilde{\lambda}_{\rm eff}(\mu) \sim \mu^{\tilde{d}-4}$ is the renormalized effective coupling at the scale μ. The vacuum energy vanishes for $\lambda = 0$ by SUSY, so $f(0) = 0$. Expanding in powers of $\tilde{\lambda}$ therefore gives (using $\Lambda_{\rm IR} \sim \omega$)

$$V_{\rm eff} \sim \omega^{\tilde{d}} + \omega^{2\tilde{d}-4} + \cdots \,, \tag{2.61}$$

which agrees with the 5D result found above.

We see that we can obtain a potential of the form eq. (2.14) with the addition of two hypermultiplets with mass parameter $c > 5/2$.[10] To obtain a large hierarchy, we need the masses of the two hypermultiplets to be approximately equal, but the tuning required is only logarithmic.

[10] Alternatively, we could use one one hypermultiplet with $c \simeq 9/2$ (corresponding to $\tilde{d} \simeq 6$) together with Casimir energy.

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2.5 SUSY breaking from stabilization

We now consider the size of SUSY breaking on the IR brane arising from the radius stabilization dynamics.

The hypermultiplet F terms can give rise to direct SUSY breaking from couplings on the IR brane of the form (in units where $\Lambda_{\rm UV} = 1$)

$$\begin{aligned}\Delta\mathcal{L}_{\rm IR} \sim & \int d^4\theta\, \Phi^\dagger\Phi \left(Q^\dagger Q + H^\dagger H\right) + \left(\int d^2\theta\, \Phi W^\alpha W_\alpha + \text{h.c.}\right) + \\ & + \int d^2\theta\, \Phi Q^2 H + \text{h.c.} + \\ & + \int d^4\theta \left(\Phi^\dagger H^2 + \Phi^\dagger \Phi H^2 + \text{h.c.}\right), \end{aligned} \tag{2.62}$$

where Q and H are matter and Higgs fields localized on the IR brane, and W_α is the field strength for standard model gauge fields, also localized on the IR brane. Just as in 'minimal SUGRA' models, this gives rise to scalar masses, gaugino masses, A terms, as well as μ and $B\mu$ terms naturally of the same size, namely

$$M_{\rm SUSY} \sim F_{\rm IR} \sim \omega^d \sim \Lambda_{\rm IR}\omega^{\tilde{d}-4}, \tag{2.63}$$

where we have restored the mass scales in the last step. Note that $M_{\rm SUSY} \ll \Lambda_{\rm IR}$, as expected. As a mechanism for SUSY breaking, this has the attraction of simplicity and elegance; in particular, it generates μ and $B\mu$ terms without additional complicated structure [29]. There is however no explanation of the absence of squark mixing required to avoid large flavor-changing neutral currents. This type of SUSY breaking therefore requires additional flavor structure at high scales, such as the models of refs. [30].

In the 4D CFT interpretation, the effects parameterized by eq. (2.62) represent SUSY breaking effects of the composite CFT states from irrelevant CFT operators. Like the operators in the 5D description, these are very generic effects that are expected to be present unless there are special symmetries that forbid them. It is remarkable that these effects can naturally give rise to all required SUSY breaking at the same scale.

Another potentially important source of SUSY breaking is the radion F term generated from the stabilization dynamics. The stabilization breaks SUSY, as can be seen from the nonzero value of the hypermultiplet F terms in the bulk. To compute the radion F term, we will use the technique of analytic continuation into superspace. The strategy is to write all SUSY breaking terms in the lagrangian using superfield spurions, and obtain a superfield form of the 4D effective potential. We can then find the dependence on the radion F term in the 4D effective theory by promoting ω to a chiral superfield [19].

We begin by writing the SUSY-breaking potential on the UV brane in terms of a superfield spurion:

$$\Delta\mathcal{L}_5 = \delta(y) \int d^4\theta\, S(\Phi, \Delta\Phi), \tag{2.64}$$

where

$$S(\Phi, F) = -\theta^4\, U(\Phi, F), \tag{2.65}$$

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and we use the abbreviations

$$\Delta = -\frac{1}{4}D^2\,, \qquad \bar{\Delta} = -\frac{1}{4}\bar{D}^2\,. \tag{2.66}$$

The $\tilde{\Phi}$ and Φ superfield equations of motion are

$$\begin{aligned}
&e^{-2\sigma}\bar{\Delta}\tilde{\Phi}^\dagger + e^{-3\sigma}\left[\partial_y\Phi + \left(c-\frac{3}{2}\right)\sigma'\Phi\right] = 0\,,\\
&e^{-2\sigma}\bar{\Delta}\Phi^\dagger + e^{-3\sigma}\left[-\partial_y\tilde{\Phi} + \left(c+\frac{3}{2}\right)\sigma'\tilde{\Phi}\right] =\\
&\qquad = -\delta(y)\bar{\Delta}\left[\frac{\partial S}{\partial\Phi} + \Delta\left(\frac{\partial S}{\partial F}\right)\right] - \delta(y-\ell)\omega^3\frac{\partial W}{\partial\Phi}\,.
\end{aligned} \tag{2.67}$$

The solutions are

$$\Phi = \Phi_0 e^{-(c-3/2)\sigma} - \frac{1}{(2c+1)k}\bar{\Delta}\tilde{\Phi}_0^\dagger e^{(c+5/2)\sigma}\,, \tag{2.68}$$

$$\tilde{\Phi} = \frac{\sigma'}{k}\left[\tilde{\Phi}_0 e^{(c+3/2)\sigma} - \frac{1}{(2c-1)k}\bar{\Delta}\Phi_0^\dagger e^{-(c-5/2)\sigma}\right], \tag{2.69}$$

where Φ_0 and $\tilde{\Phi}_0$ are superfield constants of integration, determined by the jump equations on the UV and IR branes:

$$\begin{aligned}
\tilde{\Phi}_0 - \frac{1}{(2c-1)k}\bar{\Delta}\Phi_0^\dagger &= \frac{1}{2}\bar{\Delta}\left[\frac{\partial S}{\partial\Phi} + \Delta\left(\frac{\partial S}{\partial F}\right)\right]_{\rm UV}\,,\\
\tilde{\Phi}_0\omega^{-(c-3/2)} - \frac{1}{(2c-1)k}\bar{\Delta}\Phi_0^\dagger\omega^{c-5/2} &= -\frac{1}{2}\left[\frac{\partial W}{\partial\Phi}\right]_{\rm IR}\,.
\end{aligned} \tag{2.70}$$

The effective potential is

$$\begin{aligned}
V_{\rm eff} = &\int d^4\theta\left[-S + \frac{1}{2}\left(\Phi\frac{\partial S}{\partial\Phi} + F\frac{\partial S}{\partial F} + \text{h.c.}\right)\right]_{\rm UV} +\\
&+ \int d^2\theta\,\omega^3\left[-W + \frac{1}{2}\Phi\frac{\partial W}{\partial\Phi}\right]_{\rm IR} + \text{h.c.}
\end{aligned} \tag{2.71}$$

Note that this depends on ω implicitly via the solutions for Φ_0 and $\tilde{\Phi}_0$. It is easily verified that these expressions reproduce the component expressions given earlier.

The expression eq. (2.71) can be directly analytically continued into superspace by promoting ω to a chiral superfield. To compute the VEV of F_ω, we need the linear term in F_ω in the effective potential. (The leading quadratic term in F_ω comes from eq. (2.7).) Note that there is no contribution to the coefficient of F_ω from the first term in eq. (2.71), which is a function of UV quantities. Even though this term depends implicitly on ω through Φ_0, all of the θ integrations must act on the explicit θ^4 in S to get a nonzero result. This immediately guarantees that $F_\omega/\omega \lesssim \Lambda_{\rm IR}$, and agrees with the expectation that physics associated with the UV brane does not generate an F term for the radion. A similar argument shows that the Casimir energy from bulk SUGRA does not generate a coefficient for F_ω.

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There is a nonzero linear term for F_ω from the second term in eq. (2.71), which depends on IR quantities. This is

$$V_{\text{eff}} = \omega^2 \Bigg[3\left(-W + \frac{1}{2}\Phi\frac{\partial W}{\partial \Phi}\right) + \\ + \frac{1}{2}\omega\frac{\partial \Phi_{\text{IR}}}{\partial \omega}\left(-\frac{\partial W}{\partial \Phi} + \Phi\frac{\partial^2 W}{\partial \Phi^2}\right)\Bigg]_{\text{IR}} F_\omega + \text{h.c.} + \cdots . \tag{2.72}$$

Note that the coefficient of F_ω vanishes identically if $W = \lambda\Phi^2$. This makes sense, since this term preserves conformal symmetry (λ is dimensionless).

For the example considered above, $W = \kappa\Phi$ and $\Phi_{\text{IR}} \sim \omega^{d-4}$, we have

$$\frac{F_\omega}{\omega} \sim \omega\Phi_{\text{IR}} \sim \Lambda_{\text{IR}}\omega^{d-4} . \tag{2.73}$$

The left-hand sideis the order parameter for anomaly mediated SUSY breaking (AMSB) on the IR brane [8]. Comparing to eq. (2.63), we see that AMSB from this source is always smaller than direct SUSY breaking, even without taking into account the loop suppression factors in AMSB.

Integrating out F_ω also generates SUGRA corrections to the radion potential. This gives $\Delta V_{\text{eff}} \sim \omega^{2d-4}$, which is negligible compared to the $\omega^{\tilde{d}}$ term found above.

We can obtain an additional contribution to the VEV of F_ω/ω if there is a small constant term in the superpotential on the IR brane:

$$\Delta W = C \tag{2.74}$$

gives

$$\frac{F_\omega}{\omega} \sim \Lambda_{\text{IR}}\left(\frac{C}{\Lambda_{\text{UV}}^3}\right) . \tag{2.75}$$

As discussed above, if there are no small parameters in the theory, we have $C \sim \Lambda_{\text{UV}}^3$ and the IR scale is too low. However $C \ll \Lambda_{\text{UV}}^3$ is natural because C breaks a $\text{U}(1)_R$ symmetry. In the 4D CFT interpretation, the CFT has a small parameter that breaks the $\text{U}(1)_R$ symmetry and dynamically generates a small F term for the dilaton. In this scenario, SUSY breaking can be dominated by AMSB. This is attractive because it automatically gives flavor-blind SUSY breaking.

3. Phenomenology and cosmology

In this section, we make some brief remarks about phenomenology and cosmology.

First we consider the radion, which is the only model-independent new degree of freedom in this framework. The radion is localized near the IR brane, and so its couplings to visible matter are suppressed by powers of Λ_{IR} rather than being Planck suppressed. The corresponding mode in 4D CFT language is the dilaton, which is a bound state of the CFT dynamics at the scale $\Lambda_{\rm IR}$. This modulus therefore couples more strongly than gravitational moduli, making the cosmology much safer. Because the radion decouples in the conformal limit, its couplings will be suppressed by loop factors, e.g.

$$\Delta\mathcal{L} \sim \frac{g^2}{16\pi^2}\frac{\hat{\omega}}{\Lambda_{\rm IR}}F^{\mu\nu}F_{\mu\nu}\,, \tag{3.1}$$

where $\hat{\omega}$ is the canonically normalized radion field and $F_{\mu\nu}$ is a standard model field strength. The radion mass is given by eq. (2.16) with $n_1 = \tilde{d}$:

$$m_{\rm radion} \sim \Lambda_{\rm IR}\omega^{\tilde{d}-4}\,. \tag{3.2}$$

For the large values of $\Lambda_{\rm IR}$ we are considering, the radion effectively decouples in collider experiments.

There are two scenarios for SUSY breaking that we will discuss. We first consider the case where SUSY is broken in the visible sector by direct mediation from the CFT. In the 5D description, the SUSY breaking effects arise from the couplings in eq. (2.62). In this case, we have (see eq. (2.63))

$$\omega \sim (10^{-16})^{1/(\tilde{d}-3)}\,. \tag{3.3}$$

For example, for $\tilde{d} = 5$, we have $\omega \sim 10^{-8}$, which gives $\Lambda_{\rm IR} \sim 10^{10}\,\text{GeV}$. From eqs. (3.1) and (2.63), we see that the radion mass is of order 100 GeV, for any value of $\tilde{d}$. There is no model-independent prediction for the pattern of soft masses in this scenario. If we make the plausible assumption that the scalar masses and A terms are universal at the fundamental scale, the SUSY breaking pattern is the same as in 'minimal SUGRA.'

We now briefly consider radion cosmology in this scenario. We expect radion oscillations to dominate the universe when the temperature drops below the radion mass of order 100 GeV. The reheat temperature is of order

$$T_{\rm RH} \sim 3\,\text{TeV}\left(\frac{\Lambda_{\rm IR}}{10^{10}\,\text{GeV}}\right)^{-1}. \tag{3.4}$$

This is easily large enough for a realistic cosmology.

The other scenario we discuss is that SUSY breaking is dominated by a nonzero constant superpotential. In this case, $m_{\rm radion} \gg 100\,\text{GeV}$ and SUSY breaking in the observable sector is anomaly-mediated. This also requires additional structure to obtain a realistic superpartner mass spectrum, but the mechanisms of e.g. refs. [6, 7] can be used to obtain realistic models. The detailed phenomenology is model-dependent in this case also. Radion cosmology is easily realistic. For example, for $\tilde{d} = 5$ and $\omega = 10^{-7}$, we have $\Lambda_{\rm IR} \sim 10^{11}\,\text{GeV}$, $m_{\rm radion} \sim 10\,\text{TeV}$, and $T_{\rm RH} \sim 10^5\,\text{GeV}$.

4. Conclusions

We have proposed a new paradigm for SUSY and SUSY breaking. In this approach, fundamental physics is completely non-supersymmetric, but the theory flows toward a supersymmetric conformal fixed point at low energies. SUSY is therefore an accidental

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approximate symmetry, and SUSY breaking at low energies is explicit rather than spontaneous. Remarkably, many of the features required of a realistic SUSY model follow very generically from the property that the fixed point is attractive, i.e. there are no relevant SUSY breaking perturbations of the fixed point. SUSY breaking at low energies naturally arises because the approach to the fixed point is halted due to a potential generated by irrelevant operators. This is contrast to the standard paradigm of spontaneous SUSY breaking which requires carefully chosen SUSY breaking sectors that are generally not robust against perturbations. Also, in the present framework all required SUSY breaking in the visible sector naturally occurs at the same scale with no small input parameters. If there is a small parameter due to an approximate broken $\mathrm{U}(1)_R$ symmetry, SUSY breaking can be dominated by anomaly-mediated SUSY breaking.

The detailed phenomenology of visible sector SUSY breaking is model-dependent, but there are some general consequences of this framework. Because SUSY is not an exact symmetry in the UV, there is no Goldstino. Also, the gravitational sector is completely non-supersymmetric. This means that there is no gravitino, and hence no problems with gravitino cosmology. Also, we expect that there are no gravitational moduli, which cause grave cosmological difficulties in e.g. string theory. There is a dilaton modulus in the CFT, but it couples much stronger than gravity, and does not give rise to cosmological difficulties. Finally, this framework requires that the standard model is composite at a scale

$$\Lambda_{\mathrm{IR}} \lesssim 10^{11}\,\mathrm{GeV}\,. \tag{4.1}$$

This is below the unification scale, but one-step gauge coupling unification can naturally be explained by accelerated unification [13].

We conclude that this framework is an attractive alternative to the standard paradigm of spontaneously broken supersymmetry. We have used results from 5D models to illustrate these points, but we have argued that they are very general.

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