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SUPERGRAVITY WITH AN R-SYMMETRY

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# TUNING THE COSMOLOGICAL CONSTANT IN N = 1 SUPERGRAVITY WITH AN R-SYMMETRY*

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## Abstract

An explicit example of a hidden sector for the spontaneous breaking of N = 1 local supersymmetry with an R-symmetry is constructed. It is shown that the cosmological constant of this example can be made to vanish by tuning the parameters of a non-trivial Kähler potential.

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Models of particle physics based upon the coupling of supersymmetric gauge theories to $N = 1$ supergravity [1,2,3] have recently been considered by several authors [4,5]. In these models, supersymmetry is spontaneously broken by the vacuum expectation values (VEVs) of a set of chiral fields $z_H$ which form a "hidden sector"[2,5]. The remaining chiral fields $z_O$ form an "observable sector" which includes ordinary matter and Higgs fields and which would have unbroken global supersymmetry in the absence of the coupling to supergravity. In particular, it is assumed that the superpotential takes the form

$$f = f_H(z_H) + f_O(z_O), \tag{1}$$

in which $f_H$ and $f_O$ depend only on the chiral fields of the hidden and observable sectors, respectively.

A superpotential of the form (1) can be made strongly natural in a non-renormalizable theory only if the theory possesses an R-symmetry [6]. Unfortunately, the simplest explicit example of a hidden sector [7] is incompatible with the presence of an R-symmetry. This is because the hidden superpotential $f_H$ of that example contains an additive constant which must be chosen to make the cosmological constant vanish. In this paper, we construct an explicit example of a hidden sector which possesses an R-symmetry. In this example, it is possible to make the cosmological constant vanish by adjusting the parameters of a non-trivial Kähler potential.

The two sector hypothesis of Eq. (1) provides a convenient mechanism for inserting a realistic pattern of supersymmetry breaking into models coupled to $N = 1$ supergravity. Although the supersymmetry breaking VEVs of the hidden sector in such models are typically of order the Planck mass $M_P$, one may choose parameters to make the effects of supersymmetry breaking on the observable fields and gauge fields be characterized by a gravitino mass $m_g$ of order 100-1000 GeV. This is possible when the hidden fields are neutral under all gauge symmetries and when the two sectors are separated in a superpotential of the form given by Eq. (1). Under these assumptions the two sectors communicate only through non-renormalizable interactions which are mediated by the auxilliary fields of the supergravity multiplet and which are supressed at low energies by powers of the Planck mass [5]. At energies below $M_P$, the hidden fields can be integrated out, and the dynamics of the observable sector can be described by an effective theory in which global supersymmetry is explicitly broken by soft operators at a scale of order $m_g$ [2–5].

Although terms in the superpotential are not generated by radiative corrections, a strongly natural superpotential may be desired on esthetic

grounds. The separation of the two sectors in the superpotential (1) cannot be enforced by any ordinary symmetries. Indeed, since both $f_H$ and $f_O$ must be invariant under these symmetries, terms such as $f_H f_O$ through which the sectors interact are allowed. However, a superpotential of this form can be made strongly natural if the theory possesses an R-symmetry (under which all terms of the superpotential transform with the same non-trivial phase factor) provided that the chiral fields of the theory all have positive R-characters.

The presence of a continuous R-symmetry is phenomenologically acceptable provided that it is spontaneously broken at a sufficiently large scale. In particular, if the R-symmetry were broken by the VEVs of the hidden fields, the resulting Goldstone boson would decouple below the Planck scale. The R-symmetry would then appear to be explicitly violated in the low energy effective theory by the soft operators which break supersymmetry, and phenomenologically acceptable gauge fermion masses could then be generated. Alternatively, a discrete R-symmetry could be used to forbid superpotential interactions between the sectors up to terms of sufficiently high order in the observable fields.

Lagrangians for the coupling of chiral fields to $N = 1$ supergravity [1,3] can be written in terms of a single function $G(z^*_\alpha, z^\alpha)$ of the chiral fields $z^\alpha$ and their complex conjugates. This function is customarily divided into a superpotential $f(z^\alpha)$ which depends only on the chiral fields and a Kähler potential $K(z^*_\alpha, z^\alpha)$ which depends on both the $z^\alpha$ and their conjugates:

$$G(z^*_\alpha, z^\alpha) = K(z^*_\alpha, z^\alpha) + \ln|f(z^\alpha)|^2. \tag{2}$$

In particular, we will make use of the form of the scalar potential (in units with $\kappa = 8\pi/M_P^2 = 1$)

$$\begin{aligned} V &= e^G\left\{G^{,\alpha}[G^{,\beta}_{,\alpha}]^{-1}G_{,\beta} - 3\right\} \\ &= e^K\left\{(f^{*,\alpha} + K^{,\alpha}f^*)[K^{,\beta}_{,\alpha}]^{-1}(f_{,\beta} + K_{,\beta}f) - 3|f|^2\right\}, \end{aligned} \tag{3}$$

in which, for example, $G^{,\beta\cdots}_{,\alpha\cdots}$ denotes the partial derivative of $G$ with respect to $z^\alpha, \cdots$ and $z^*_{\beta,}\cdots$.

For simplicity, we confine the discussion to hidden sectors which involve only a single chiral field $z_H$. Up to a field redefinition, the requirement

of an R-invariant theory restricts the hidden superpotential to be $f_H = az_H$ in which $a$ is a constant of order $m_g M_P$. This superpotential possesses an R-symmetry under which $z_H$ has R-character two, and, in the absence of observable fields, this will be a symmetry of the theory provided that the Kähler potential $K$ depends only on $x = |z_H|^2$. The scalar potential for the hidden sector alone is then [3]

$$V_H(x) = |a|^2 e^{K(x)} U(x), \tag{4}$$

in which

$$U(x) = \frac{(1 + xK')^2}{K' + xK''} - 3x, \tag{5}$$

where primes denote differentiation with respect to $x$. For the cosmological constant to vanish in the absence of the observable sector, we must choose $K(x)$ so that:

(i) $K(x)$ is well-behaved for $x \geq 0$

and

(ii) $U(x) > 0$ for $x \geq 0$ with the exception that $U(\overline{x}) = 0$ for some $\overline{x} \geq 0$.

These conditions imply that the true vacuum of $V_H$ has $|\langle z_H \rangle|^2 = \overline{x}$ and that the vacuum energy vanishes.

It is interesting to note that condition (ii) is trivial to satisfy. Indeed, we may choose $U(x)$ to be any function which satisfies this condition and regard Eq. (5) as a differential equation for the Kähler potential. We then find that

$$K(x) = -\int_{x_0}^{x} \frac{dy}{y}[1 + W(y)], \tag{6}$$

where

$$W(x) = \left[\int_{x_1}^{x} \frac{dy}{3y + U(y)}\right]^{-1}, \tag{7}$$

with $x_0$ and $x_1$ denoting arbitrary constants. Of course, we must still check that condition (i) is satisfied for this Kähler potential.

To obtain an explicit example, we choose

$$U(x) = \frac{(1 + x + \alpha x^2)^2}{1 + 2\alpha x} - 3x, \tag{8}$$

where $\alpha$ is a positive constant. For $x_0 = 0$ and $x_1 = -\infty$, this corresponds to the Kähler potential

$$K(x) = x + \frac{1}{2}\alpha x^2. \tag{9}$$

Figure 1 shows the resulting hidden sector potential $V_H(x)$ for three values of $\alpha$. When this parameter takes the special value

$$\overline{\alpha} = \frac{1}{4} + \frac{3}{16}\left[(2 + \sqrt{3})^{\frac{1}{3}} + (2 - \sqrt{3})^{\frac{1}{3}}\right], \tag{10}$$

the potential $V_H$ is positive with the exception of a double root which occurs at

$$\overline{x} = \left\{\frac{1}{2} + \frac{1}{4}\left[(7 + 4\sqrt{3})^{\frac{1}{3}} + (7 - 4\sqrt{3})^{\frac{1}{3}}\right]\right\}^{-1}. \tag{11}$$

The true vacuum in the absence of any observable sector then has $|\langle z_H \rangle|^2 = \overline{x}$ and zero cosmological constant.

To conclude the paper, we briefly consider the incorporation of an R-invariant observable sector into theories with the hidden sector superpotential $f_H = a z_H$. For simplicity, we ignore gauge interactions and assume that the Kähler potential is invariant under $U(N)$ rotations of the N chiral fields $z_H$ and $z^a$, where the $z^a$ are the scalar fields of the observable sector — i.e., we take $K$ to be a function of only

$$u = |z_H|^2 + z^*_a z^a. \tag{12}$$

The scalar potential is then given by

$$V = e^K \left\{ \frac{1}{K'} |F_H|^2 + \frac{1}{K'} F_a^* F^a - L|G|^2 - 3|az_H + f_O|^2 \right\}, \tag{13}$$

where

$$F_H^* = a + K' z_H^* (az_H + f_O)$$

$$F_a^* = f_{O,a} + K' z_a^* (az_H + f_O)$$

$$G = az_H + z^a f_{O,a} + uK'(az_H + f_O)$$

and

$$L = \frac{K''}{K'(K' + uK'')}. \tag{14}$$

Also, primes denote derivatives of K with respect to its argument, while $f_{O,a}$ are the derivatives of the observable superpotential $f_O$ with respect to the fields $z^a$.

The simplest case is that in which the observable sector describes weak interaction scale physics with no grand unification being involved. In this case, the dimensionful parameters of $f_O$ have the characteristic scale $m_g$, and we may take the observable fields $z^a$ to be of order $m_g$ so that $f_O$ is of order $m_g^3$. Expanding the potential (14) in powers of $m_g$, we find that

$$V = V_H + V_O + O(m_g^5). \tag{15}$$

in which $V_H$ is the potential (4) and (5) of the hidden sector by itself, while $V_O$ consists of all terms of order $m_g^4$ which involve observable fields. It is assumed that $V_H$ is minimized with zero vacuum energy at some value $\overline{z}_H$ of $z_H$ and for some choice $\overline{\alpha}_n$ of the parameters $\alpha_n$ which appear in the Kähler potential:

$$\frac{\partial V_H}{\partial z_H}|_{\overline{z}_H, \overline{\alpha}_n} = 0 \quad \text{and} \quad V_H|_{\overline{z}_H, \overline{\alpha}_n} = 0. \tag{16}$$

To order $m_g^4$, the potential of the low energy effective theory for the observable sector is just $V_O|_{\overline{z}_H,\overline{\alpha}_n}$ plus a constant which arises from shifts of order $m_g^2$ in the values of the parameters $\alpha_n$. These shifts may be chosen to cancel the vacuum energy which is obtained from minimizing $V_O$.

The presence of the hidden sector induces soft operators which break supersymmetry in the low energy potential [2,5]. To display these operators, we write $V_O$ in terms of fields $\tilde{z}^a = \sqrt{K'(\overline{x})}\, z^a$ which have canonically normalized kinetic terms and a low energy superpotential

$$f_{LE}(\tilde{z}^a) = e^{\frac{1}{2}K(\overline{x})} f_O([K'(\overline{x})]^{-\frac{1}{2}}\tilde{z}^a), \tag{17}$$

in which $\overline{x} = |\overline{z}_H|^2$. After absorbing the phase of $a\overline{z}_H$ into $f_{LE}$, we have

$$V_O = f_{LE}^{*,a} f_{LE,a} + Am_g \cdot 2Re(f_{LE}) + Bm_g \cdot 2Re(\tilde{z}^a f_{LE,a}) + Cm_g^2 \tilde{z}_a^* \tilde{z}^a \tag{18}$$

where $m_g = |a\overline{z}_H| \exp[\frac{1}{2}K(\overline{x})]$ is the gravitino mass [1] and in which the derivatives of $f_{LE}$ are with respect to $\tilde{z}^a$. The soft terms linear in $f_{LE}$ provide explicit R-symmetry breaking in the low energy theory. Also, we note that the coefficient C is not constrained by the mass-squared sum rule [3], since this does not apply in theories with a non-trivial Kähler potential. Using the conditions (16), we find that the coefficients of the soft operators are given by

$$A = \frac{-3}{1 + \overline{x}K'(\overline{x})}$$

$$B = \frac{3}{1 + \overline{x}K'(\overline{x})} - \frac{1}{\overline{x}K'(\overline{x})}$$

and

$$C = \frac{1}{[\overline{x}K'(\overline{x})]^2}. \tag{19}$$

In the explicit example (Eq. 9 with $x$ replaced by $u$ and $\alpha = \overline{\alpha}$ from Eq. 10), these coefficients are $A = -1.31$, $B = 0.53$ and $C = 0.61$.

When gauge interactions are included, supersymmetry may also be broken in the low energy theory by gauge fermion masses linear in $m_g$. In the

present context, these masses vanish at tree level, because the presence of a continuous R-symmetry under which all chiral fields have positive R-character forbids non-minimal F-terms of the form $[g_{AB}(z_H, z^a)W_A W_B]_F$ in which the $W_A$ are the chiral field strength superfields and $g_{AB}$ is a function of the chiral fields. However, gauge fermions masses linear in $Am_g$ or $Bm_g$ can be generated by radiative corrections at one loop.

In models incorporating grand unification, the characteristic scale for the observable sector is the unification mass $M_G$. Since this is much larger than $m_g$, a more complicated analysis, such as that of Ref. 5, is required to determine the coefficients of the soft operators. The starting point for such an analysis is the assumption that the potential in the absence of the hidden sector has a supersymmetric minimum with zero vacuum energy. That is, we assume that there is a set of VEVs $\overline{z}^a$ for the observable fields for which $f_{O,a}(\overline{z}^a) = 0$ and $f_O(\overline{z}^a) = 0$. Here we note that the second of these conditions is non-trivial in R-symmetric theories, since we cannot transfer an additive constant order $M_G^3$ from $f_O$ to $f_H$ and still take $f_H$ to be of order $m_g M_P^2$. However, this condition may be automatically satisfied in a large class of R-symmetric theories. When this condition is met, it will be possible to cancel the vacuum energy induced by the presence of an R-invariant hidden sector with $f_H = az_H$ by tuning the parameters of a non-trivial Kähler potential. The generic form of the soft supersymmetry breaking operators which emerge in grand unified theories of this type is given in Ref. 5.

In this paper, we have constructed a simple example to demonstrate the possibility of cancelling the cosmological constant in $N = 1$ supergravity models which possess an R-symmetry. This was motivated by the attractiveness of using an R-symmetry to enforce the separation of the hidden and observable sectors in the superpotential. Although the form of the Kähler potential in this example was not natural, we are not particularly interested here in the form of the Kähler potential, so long as there is sufficient freedom to adjust the parameters to cancel the vacuum energy. Of course, this adjustment is not natural and must be repeated in each order of perturbation theory.

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# Figure Captions

Figure 1. The hidden sector potential defined by equations (4), (5) and (9) for the three values of the parameter $\alpha$: (A) $\alpha = 0$, (B) $\alpha = \overline{\alpha}$ from Eq. (10), and (C) $\alpha = 1.5$.

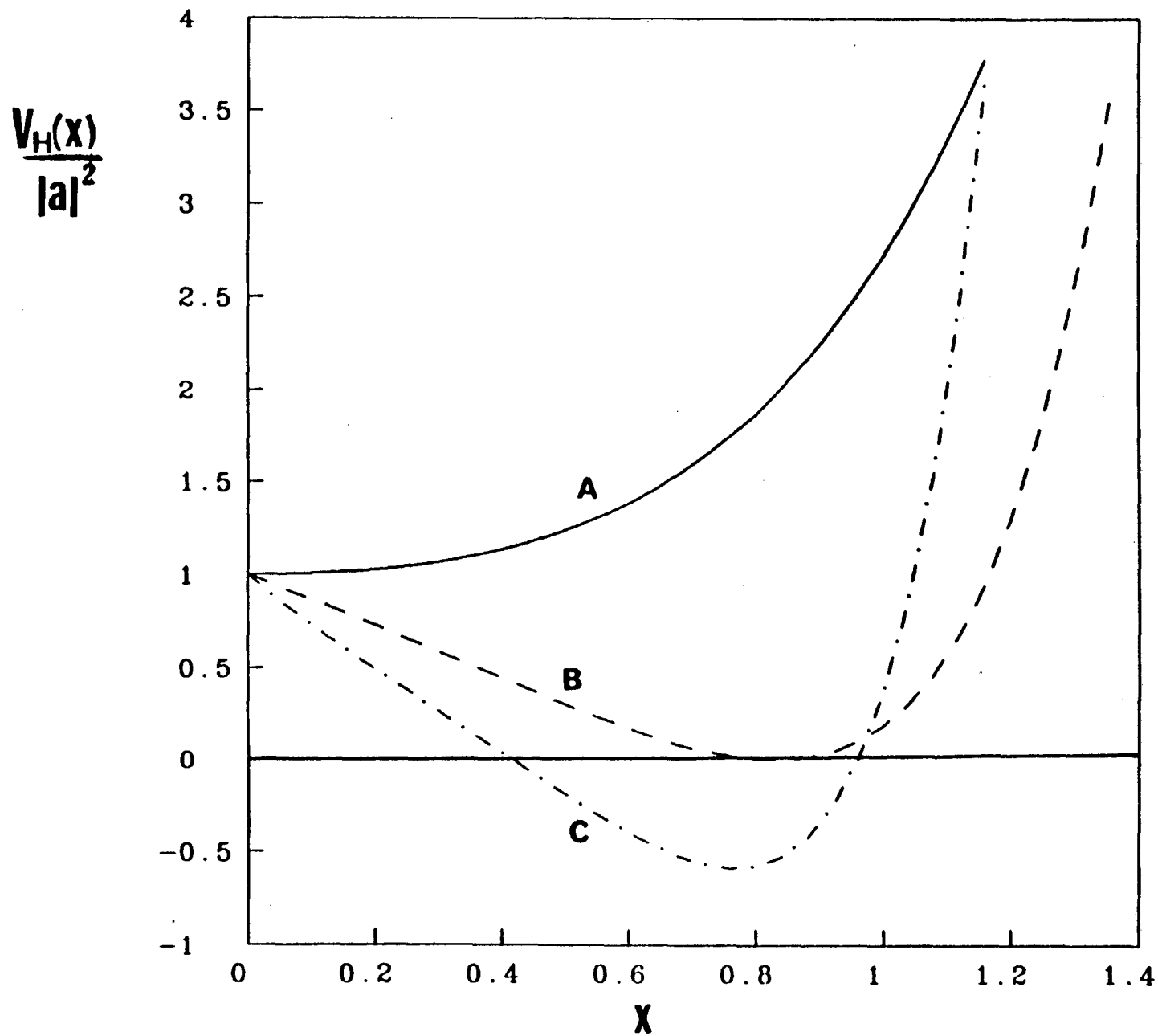


Figure 1

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