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NON-COMPACT SUPERGRAVITY SOLVES PROBLEMS

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ABSTRACT

We present a class of no-scale supergravity models based on maximally symmetric non-compact Kähler manifolds which

- have $m_{3/2} \ll 0(m_W)$;
- respect all particle physics and cosmological constraints on gravitinos;
- allow the QCD θ parameter to relax harmlessly to zero;
- contain an axion-like particle compatible with particle physics and astrophysics constraints;
- have a hidden sector that does not contain excessive axion vacuum energy.

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If one is to generate [1-4] a weak gauge hierarchy, $m_W = m_P \exp [-O(1)/\alpha]$, in the context of supersymmetry, one needs a supergravity theory with a flat potential [5] so that the scale of supersymmetry breaking is undetermined at the tree level and may be fixed by radiative corrections. Such flat potentials appear in 'no-scale' supergravity models based on non-compact Kähler manifolds, the maximally symmetric coset spaces $SU(n,1)/SU(n) \times U(1)$ [2]. Phenomenologically acceptable no-scale models with $m_{3/2} \gg m_W$ [3], $m_{3/2} = O(m_W)$ [1,2], or $m_{3/2} \ll m_W$ [4] have been proposed. We have previously pointed out [4] that such models offer a novel solution to the strong CP problem [6], but our previous paper did not explore all the particle physics, astrophysical, and cosmological constraints on such a model.

In this paper we discuss in more detail our proposed [4] solution to the strong CP problem, which involves dynamical relaxation of the θ vacuum parameter. We define the range of parameters in no-scale models for which this vacuum relaxation is compatible with a wide range of particle physics constraints such as the absence [7] of an unobserved axion-like particle and the weakness [8] of gravitino couplings to ordinary matter, of astrophysical constraints [9,10] on such an 'axion', and cosmological constraints on vacuum energy and entropy generation in the hidden sector [11] as well as gravitino regeneration [12]. We also mention an interesting class of models [13] with $m_{\tilde{\gamma}} > m_{3/2}$ and both masses $O(10)$ GeV, in which the sparticle spectrum can be 'light' as favoured by one possible interpretation [14] of recent CERN $p\bar{p}$ Collider data [15].

Motivated by the desire to generate a weak gauge hierarchy, $m_W = m_P \exp [-O(1)/\alpha]$, we work in the context of 'no-scale' supergravity models with an underlying non-compact $SU(n,1)/SU(n) \times U(1)$ manifold structure [2]. For illustrative simplicity, we present here examples based on a $SU(1,1)/U(1)$ submanifold with a single complex field z parametrizing the hidden sector:

$$G = -3 \ln(z + z^*) + (\text{terms involving observable fields } y) . \quad (1)$$

The theory requires [16] for its specification, in addition to the real Kähler potential (1), a chiral function $h_{ab}(z, y)$ which fixes the gauge kinetic terms related to the θ vacuum parameter and gaugino masses. In general, h_{ab} transforms as the

symmetric product of two adjoint representations of the gauge group. To keep things simple, we examine the choice $h_{ab}(z,y) = \delta_{ab} h(z)$ in which case the relevant $N = 1$ supergravitational couplings of the Yang-Mills smultiplet are [16]

$$\begin{aligned} e^{-1} \mathcal{L} = & -\frac{1}{4} \text{Re } h(z) [F_{\mu\nu}^a F^{\mu\nu a} + \bar{\lambda}^a \not{D} \lambda^a] \\ & + \frac{i}{4} \text{Im } h(z) [\epsilon^{\mu\nu\rho\lambda} F_{\mu\nu}^a F_{\rho\lambda}^a - \frac{1}{2} e^{-1} D_\mu (e \bar{\lambda}^a \gamma_5 \gamma^\mu \lambda^a)] \\ & + \frac{1}{4} e^{G/2} G_z (G_z z^*)^{-1} h_{z^*}^*(z^*) \bar{\lambda}_R^a \lambda_R^a + \dots \end{aligned} \quad (2)$$

where the z and z^* subscripts denote derivatives. After the correct rescaling of the gaugino fields, their SUSY-breaking masses are given by

$$m_{\tilde{V}} = \left| m_{3/2}^{1/3} \left[\text{Im } \text{Re } h^* \right]_{z^*} \right| \quad (3)$$

Note that the usual proportionality $m_{\tilde{V}} \propto m_{3/2}$ is not apparent [Eq. (3)] in these $SU(1,1)/U(1)$ models. Indeed, with an appropriate choice of $h(z)$, any desired ratio between $m_{\tilde{V}}$ and $m_{3/2}$ may be obtained. For example, in a previous paper [4] we proposed consideration of

$$h(z) = e^{iA z^2} + iC \quad (4)$$

and we now explore in more detail the properties of this ansatz.

The single complex field z in our hidden sector (1) contains two real components which we may write as $z = u + iv = \rho e^{i\phi}$. It is a general feature of no-scale models [1-4] that $m_{3/2}$ is specified by the real part of z ,

$$m_{3/2} = (2u)^{-3/2} \quad (5)$$

which can be dynamically determined by a mechanism of dimensional transmutation associated with weak radiative corrections [1]. The second degree of freedom in z remains to be fixed, and we propose [4] that it is specified by the dynamical relaxation of the QCD θ vacuum parameter associated with non-perturbative QCD effects. To see how the two components of z can be specified, we consider the sample model [Eqs. (1) to (4)], with the choice $c = 0$ in expression (4). In models of this type, the weak radiative corrections fix directly the scale of global SUSY breaking characterized by the gaugino mass:

$$m_{\tilde{V}} = qA [1 + \theta^2]^{\frac{1}{2}} \rho^{q-1} m_{3/2}^{\frac{1}{3}} \quad (6)$$

so that $m_{\tilde{V}} = O(100)$ GeV, while $m_{3/2}$ may be very different. In Eq. (6), θ is the QCD vacuum angle parameter, which is seen from the first and second terms of Eq. (2) to have the value

$$\theta = \tan A \operatorname{Re}(z^q) = \tan(A \rho^q \cos q\phi) \quad (7)$$

Because z is a field variable, θ is a dynamical quantity which can be fixed by minimizing the energy of the system. Non-perturbative QCD effects [6] give

$$V = V_0 + \frac{1}{2} \left(\frac{d^2 V}{d\theta^2} \right)_{\theta=0} \theta^2 + \dots \quad (8)$$

where [17]

$$\left(\frac{d^2 V}{d\theta^2} \right)_{\theta=0} = \frac{f_{\pi}^2 m_{\pi'}^2}{4N_f} \quad (9)$$

with $N_f = 3$. Therefore, θ can relax to zero so that

$$\rho^q \cos q\phi = 0 \pmod{\frac{\pi}{A}} \quad (10)$$

Knowing that $\theta = 0$, we can now rewrite Eq. (6) in the form

$$m_{3/2} = \left(\frac{2 \cos \phi}{qA} m_{\tilde{V}} \right)^p : p = \frac{1}{1 - \frac{2}{3}q} \quad (11)$$

or in the form

$$\rho^{q-\frac{3}{2}} (\cos \phi)^{-\frac{1}{2}} = \frac{\sqrt{2} m_{\tilde{V}}}{qA} \quad (12)$$

Equations (10) and (12) clearly suffice to determine, up to a discrete ambiguity, ρ and ϕ , i.e. all components of the hidden sector field z , whilst Eq. (11) tells us that

$$m_{3/2} = \alpha(1) \times m_{\tilde{V}}^p \quad (13)$$

with p a model-dependent power. Conventional supergravity models [18] have assumed $m_{3/2} = O(m_{\tilde{V}})$, corresponding to $p = 1$ in our case. For us a more general choice of

$p(q)$ is possible, and we will investigate the phenomenological constraints on the value of p coming from particle physics and astrophysics.

First, however, we must estimate the masses and couplings of the physical particle degrees of freedom in the 'hidden' z sector. The weak interaction radiative corrections generate [1,19] a non-trivial dependence of the effective potential on $t \equiv \ln(m_{\tilde{V}}^2/\mu_0^2)$, where μ_0 is the renormalization scale where spontaneous symmetry breaking becomes possible:

$$V_{RC} = O(\mu_0^4) \left[-e^{2t} (t - 1.537)^2 + e^{1.079} \right] \quad (14)$$

in a typical model [1], which yields

$$V_{tt} \equiv \frac{\partial^2 V}{\partial t^2} = 5.85 \times O(m_{\tilde{V}}^4) \quad (15)$$

In the neighbourhood of the minimum of this effective potential, it can be expanded quadratically in a 'Polonyi' combination of z components:

$$\mathcal{J} = [\hat{u} \sin \omega + \hat{v} \cos \omega] : \tan \omega \equiv \cot \omega_0 - [(q-1) \sin 2\omega_0]^{-1} \quad (16)$$

$$\omega_0 = \pi/2q$$

where

$$(\hat{u}, \hat{v}) \equiv \sqrt{3} m_{3/2}^{2/3} [\text{Re}(z - z_0), \text{Im}(z - z_0)] \quad (17)$$

are correctly normalized [1,2] fluctuations in z about the value z_0 which minimizes the effective potential (14). Near this minimum

$$V_{RC} \approx \frac{1}{2} m_{\mathcal{J}}^2 \mathcal{J}^2 + \dots \quad (18)$$

where

$$m_{\mathcal{J}}^2 = \left(\frac{V_{tt}}{3 m_{\tilde{V}}^4} \right) \mathcal{J}^{1-\frac{2}{p}} \left[3 + 4(q-1)(q-2) \right] m_{\tilde{V}}^4 : \mathcal{J} \equiv (2 \cos \phi)^{\frac{4p}{3-p}} \quad (19)$$

Evidently

$$m_{\mathcal{J}} = O(m_{\tilde{V}}^2) \quad (20)$$

as one would expect for a Polonyi field. However, we are not yet finished with

the effective potential for the hidden sector, which also gets contributions from non-perturbative QCD effects [6]. Consider the 'axion' combination of z components:

$$a \equiv \hat{u} \cos \alpha + \hat{v} \sin \alpha : \quad \alpha \equiv \frac{3-p}{3(p-1)} \frac{\pi}{2} \quad (21)$$

which, thanks to the second term in Eq. (2), has a coupling to gauge boson pairs,

$$\mathcal{L}_{aFF} \equiv \frac{1}{\sqrt{3}} \left(\frac{m_V^p}{m_{3/2}} \right) \frac{a}{m_A} \epsilon^{\mu\nu\rho\lambda} F_{\mu\nu}^a F_{\rho\lambda}^a \quad (22)$$

where

$$m_A \equiv m_{\tilde{V}}^{p-1} \quad (23)$$

and $(m_V^p/m_{3/2})$ is seen from Eq. (13) to be $O(1)$, whilst the sum over a in Eq. (22) runs over $SU(3)$, $SU(2)$, and $U(1)$ gauge bosons. Because of the gluon pair coupling in Eq. (22), the non-perturbative QCD effects (8) generate a term in the effective potential proportional to a^2 :

$$V_{NP}^a \approx \left(\frac{f_{\pi}^2 m_{3/2}^2}{12} \right) \left(\frac{1}{\sqrt{3}} \frac{m_V^p}{m_{3/2}} \right)^2 \left(\frac{a}{m_A} \right)^2 + \dots \quad (24)$$

Needless to say, the photon pair coupling in Eq. (22) would be responsible for 'axion' decay. We are still not quite through with the effective potential for the hidden sector, which contains additional non-perturbative QCD terms thanks to the coupling

$$\mathcal{L}_{OFF} \equiv \frac{1}{\sqrt{3}} \left(\frac{m_V^p}{m_{3/2}} \right) \frac{\sigma}{m_A} F_{\mu\nu}^a F^{\mu\nu a} \quad (25)$$

of the combination

$$\sigma \equiv \hat{u} \sin \alpha - \hat{v} \cos \alpha \quad (26)$$

of z components orthogonal to the axion. The gluon pair portion of Eq. (25) yields a dependence of the effective potential on this combination of z components:

$$V_{NP}^{\sigma} = \sum_{n=1}^{\infty} \frac{1}{n!} G_{FF}^n \left[\frac{\text{Re} h \left(\frac{1}{\sqrt{3}} \frac{m_{\tilde{V}}^p}{m_{\tilde{Z}_2}} \frac{\sigma}{m_A}, \frac{1}{\sqrt{3}} \frac{m_{\tilde{V}}^p}{m_{\tilde{Z}_2}} \frac{a}{m_A} = 0 \right)}{\text{Re} h(0,0)} \right]^n \quad (27)$$

where the G_{FF}^n are the Green functions for $n F_{\mu\nu}^a F^{\mu\nu a}$ operators at zero momentum transfer. Generally, one expects

$$G_{FF}^n = c_n \Lambda_{QCD}^4 : c_n = O(1) \quad (28)$$

and estimates of $G_{FF}^1 = \langle 0 | F_{\mu\nu}^a F^{\mu\nu a} | 0 \rangle$, G_{FF}^2 , etc., are available [20]. Thus we have

$$V_{NP}^{\sigma} = \left(\frac{1}{\sqrt{3}} \frac{m_{\tilde{V}}^p}{m_{\tilde{Z}_2}} \right) (c_1 \Lambda^4) \frac{\sigma}{m_A} + \frac{1}{2} \left(\frac{1}{\sqrt{3}} \frac{m_{\tilde{V}}^p}{m_{\tilde{Z}_2}} \right)^2 (c_2 \Lambda^4) \left(\frac{\sigma}{m_A} \right)^2 + \dots \quad (29)$$

The full effective potential for the hidden sector is

$$V = V_{RC} + V_{NP}^a + V_{NP}^{\sigma} \quad (30)$$

and we must determine its absolute minimum [shifted from z_0 by the linear term in Eq. (29)] and diagonalize the mass matrix of perturbations about this minimum in order to determine the physical mass eigenstates and their couplings.

Consider first what happens if $\Lambda^4/m_A^2 \gg m_{\tilde{V}}^4$, so that V_{RC} is negligible compared with V_{NP} in Eq. (30). Then it is easy to see from Eq. (29) that the minimum of the potential occurs where $\sigma = O(m_A)$, and the weak radiative corrections no longer determine the scale of SUSY breaking as we have hypothesized. Whilst it may be possible to construct interesting models in which non-perturbative QCD effects determine the scale of SUSY breaking, our strategy [1,4] is to generate SUSY-breaking scales from weak radiative corrections, so we restrict ourselves to the case

$$m_A > \Lambda^2 / m_{\tilde{V}}^2 \quad (31)$$

when the SUSY-breaking scale is dynamically determined by weak radiative corrections [1]. Numerically, the bound (31) corresponds to:

$$m_A = O(m_{\tilde{\nu}}^{p-1}) > O\left(\frac{\Lambda^2}{m_{\tilde{\nu}}^2}\right) = O(10^{-6}) \approx 10^{12} \text{ GeV} \quad (32)$$

$$\Rightarrow p < 1.4$$

In this case the linear term in Eq. (29) acts as a small perturbation which slightly shifts the minimum away from z_0 ,

$$\hat{u}, \hat{v} = O\left(\frac{\Lambda^4}{m_{\tilde{\nu}}^4 m_A}\right) \lesssim O\left(\frac{\Lambda^2}{m_{\tilde{\nu}}^2}\right) \ll 1 \quad (33)$$

but otherwise has a negligible effect on the effective potential and will be henceforth ignored.

What is, however, important is to ensure that the σ field is strictly parallel to the Polonyi field ζ . If it were not, the combination of the weak radiative corrections with the non-trivial non-perturbative potential for the σ field would mess up the dynamical determination of $\theta = 0$. We must ensure that the effective potential has non-trivial variation in just two orthogonal directions, one of them ζ and the other the axion field a . We have verified that in the context of our illustrative ansatz (4) for $h(z)$, this happens for an everywhere dense discrete set of values of p . (There is a particularly elegant example with $p = 1$ [13].) In such a case, the mass eigenstates will be

$$\left. \begin{aligned} \zeta: \quad m_{\zeta} &= O(m_{\tilde{\nu}}^2) \\ a: \quad m_a &= O(\Lambda^2/m_A) \end{aligned} \right\} \quad (34)$$

with photon pair couplings

$$g_{\zeta FF}, g_{a FF} = O\left(\frac{1}{m_A}\right) \quad (35)$$

We are now in a position to explore systematically the phenomenological constraints on p : $m_{3/2} = m_{\tilde{\nu}}^p$, $m_A = O(m_{\tilde{\nu}}^{p-1})$, treating in a unified way the hidden sector particles ζ, a and the gravitino.

We start with particle physics constraints on the gravitino. Fayet [8] has emphasized that the couplings of the gravitino increase as $m_{3/2}$ is decreased, and that useful bounds can be obtained from negative results of 'missing energy' experiments such as $J/\psi \rightarrow \text{invisible}$ or $e^+e^- \rightarrow (\text{invisible}) + \gamma$, interpreted as upper bounds on the associated emission of gravitinos and photinos. He [8] used an upper limit on $J/\psi \rightarrow \text{invisible}$ [21] to deduce

$$m_{3/2} > 1.5 \times 10^{-8} \text{ eV} \quad \Rightarrow \quad p = \frac{\ln m_{3/2}}{\ln m_{\tilde{\gamma}}} < 2.2 \quad (36)$$

Recent data on $e^+e^- \rightarrow (\text{invisible}) + \gamma$ enable [8] one to derive

$$m_{3/2}^2 > \frac{4}{\alpha_{em}} \bar{m}_e^4 G_N \quad (37)$$

where $\bar{m}_e$ is the minimum value of the selection mass allowed by the data [22] when they are interpreted as an upper limit on $e^+e^- \rightarrow (\tilde{\gamma}\tilde{\gamma})_{m_{\tilde{\gamma}}=0} + \gamma$. Taking [22] $\bar{m}_e = 28 \text{ GeV}$, we deduce from (37) that

$$m_{3/2} > 1.5 \times 10^{-6} \text{ eV} \quad \Rightarrow \quad p < 2.0 \quad (38)$$

Note, however, that the Fayet bound (36) disappears if $m_{\tilde{\gamma}} > m_{J/\psi}$, whilst the more stringent bound (38) disappears if $m_{\tilde{\gamma}} > 0(10) \text{ GeV}$. Therefore the bound (38) should be interpreted *cum grano salis*.

There are also particle physics bounds [7] on 'axion'-like particles which amount to a lower bound,

$$m_A = O(m_{\tilde{\gamma}}^{p-1}) > O(m_w) = O(m_{\tilde{\gamma}}) \Rightarrow p < 2.0 \quad (39)$$

However, these axion bounds are eclipsed by astrophysical bounds on energy emission by red giants [9] and by neutron stars [10]. The former would [9] emit energy in the form of light axions at a rate

$$Q_a = (\text{stellar stuff}) \times \Gamma_a / m_a^3 \quad (40)$$

while for us

$$\Gamma_a / m_a^3 = \frac{O(1)}{64\pi} \frac{1}{m_A^2} \quad (41)$$

Taking [9] $Q_a < 10^2 \text{ erg g}^{-1}\text{s}^{-1}$ yields

$$\Gamma_a/m_a^3 < O(3 \times 10^{12}) \quad (42)$$

in natural units, entailing

$$m_A > O(4 \times 10^{-8}) \Rightarrow p < 1.45 \quad (43)$$

leading to

$$m_{3/2} = O(m_{\tilde{\nu}}^p) > O(10^{-6}) \text{ GeV} = O(1) \text{ keV} \quad (44)$$

An analogous but somewhat less strict upper bound on p of about 1.6 can be derived from considering energy losses from neutron stars [10].

Notice now that after all these particle physics [formulae (36), (38), (39)] and astrophysical (43) efforts, we have almost bounded p to be within the range (32) required for consistency of our mechanism for determining the weak interaction scale by radiative corrections [1]. How about consistency with conventional cosmology? Our ζ_T field will in general produce a coherent vacuum energy comparable with that of a conventional axion [23] if we identify m_A with the conventional Higgs vacuum expectation value v . It has been argued [23] that this vacuum energy is excessive, unless $v < O(10^{12}) \text{ GeV}$, or possibly 10^{15} GeV [24]. Therefore we require

$$m_A < O(10^{15}) \text{ GeV} \Rightarrow p > 1.2 \quad (45)$$

If m_A is in the range 10^{12} to 10^{15} GeV , the a field is a prime candidate for the dark matter in the Universe. We have an analogous difficulty [11] with the cosmological energy density in our 'Polonyi' field ζ , which is more severe to the extent that the Polonyi vacuum energy scale $\propto m_\zeta^2 = O(m_{\tilde{\nu}}^4/v^2)$ is larger than $m_a^2 = O(\Lambda^4/m_A^2)$. In principle, we have at our disposal a novel means of dissipating the Polonyi vacuum energy, namely the self-interactions induced by the effective potential (29), which are much larger than in a conventional Polonyi model. Multi- σ self-interactions could convert the coherent non-relativistic Polonyi field energy into relativistic quanta, which would subsequently be redshifted harmlessly away.

This mechanism was previously shown [23] to be insufficient to erode the axion field energy, nor is it sufficient in a conventional Polonyi model. We find that even though our Polonyi self-interactions are stronger than in a conventional Polonyi model, they are still not strong enough to allow us to escape from this 'Polonyi problem'. However, it may well be that an attractive variant of our p models could be found in which this problem is solved.

How about relic gravitinos? We have seen (44) that the preferred range of gravitino masses is somewhat higher than that [$< 0(1)$ keV] permitted by the relic gravitino mass density in traditional Big Bang cosmology. However, a relatively modest amount of inflation [25] (yielding less than a factor of 10^2 in entropy) would suffice to suppress the primordial gravitino density acceptably. Gravitinos can be regenerated [12] after inflation and reheating, but the regenerated number density is seen from Eq. (14) of Ref. [12] to be acceptably low for all reheating temperatures $T_R < 0(m_P)$.

We close with some comments about supergravity models with $p = 1$. We have seen [Eq. (45)] that these models could suffer from an 'axion' problem [23] as well as the well-known 'Polonyi' problem [11]. It is often said that they also suffer from a gravitino problem [26], which could only be saved by inflation [25] if the reheating temperature $T_R < 0(10^8)$ GeV [12,27]. We would like to point here that this stringent bound on the reheating temperature can be relaxed if $m_{3/2} < m_{\tilde{\gamma}}$, something which is automatic in our models with $p > 1$, but unexpected if $p = 1$. It can, however, be arranged [13] by a devious choice of SUSY-breaking gaugino masses which are not in the usual 'unified' ratio

$$m_{\tilde{g}}/\alpha_3 = m_{\tilde{w}}/\alpha_2 = m_{\tilde{b}}/\alpha_1 \quad (46)$$

In this case gravitinos are absolutely stable and photinos decay with a lifetime $0(10^7 \text{ to } 10^8)$ s. These decays are cosmologically embarrassing, but can be suppressed to an (almost) arbitrarily low acceptable level because the photinos (unlike

gravitinos) can annihilate very efficiently by sparticle exchange [28]. Then our only worry is with the relic gravitino mass density, and this is acceptably low $[\rho_{3/2} \lesssim 0(10) \rho_{\text{matter}}]$ for $m_{3/2} = 0(10)$ GeV if $T_R \lesssim 0(10^{11})$ GeV. This expansion in the acceptable range of reheating temperatures is potentially interesting, because Higgses with interactions violating baryon number are required [29] by baryon stability to have $m_H \gtrsim 0(10^{11})$ GeV and could not be produced thermally in large numbers if $T_R \lesssim 0(10^8)$ GeV as required [12,27] if $p = 1$ and $m_{3/2} > m_{\tilde{\gamma}}$. There is no corresponding difficulty in our models with $p > 1$ (46), since the reheating temperature can be much higher and baryogenetic Higgs can be generated thermally as well as non-thermally [30].

Although we need not take too seriously the mathematical form of the chiral function $h(z)$ [Eq. (2)] which we [4] have chosen [Eq. (4)], we are excited that models with $m_{3/2} = 0(m_W^p)$: $p > 1$ (particularly $p \sim 1.2$ to 1.4) can avoid all particle physics, astrophysical, and some cosmological pitfalls, and that models with $m_{3/2} < m_W$ are natural possibilities in supergravity models [1,2] based on maximally symmetric non-compact Kähler manifolds.

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