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SUPERSTRING COMPACTIFICATION AND SUPERSYMMETRY BREAKING

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A B S T R A C T

We set out conditions which enable gaugino condensation to occur in a compactified supergravity model which may be the low energy limit of a superstring theory. We then argue that the scales of compactification, of supersymmetry breaking and of gauge symmetry breaking can be dynamically determined by radiative corrections in an effective no-scale supergravity model. The resulting low energy sector may avoid the cosmological gravitino problem - because the gravitino can decay early enough as not to disturb primordial nucleosynthesis.

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The discovery¹⁾ that certain superstring theories²⁾ are anomaly-free and may be finite has rekindled hopes that one of them may be the ultimate theory of everything (TOE), and has stimulated new attempts³⁾⁻⁹⁾ to connect superstring theories to phenomenology. Such attempts involve taking the low energy field theory limit, which is an $N = 1$, $d = 10$ supergravity theory with an $O(32)$ or more likely an $E_8 \times E_8$ gauge group¹⁰⁾, compactifying the six extra dimensions and extracting a low energy effective $N = 1$, $d = 4$ supergravity theory⁴⁾, which turns out⁵⁾ to have a "no-scale" $SU(n,1)/SU(n) \times U(1)$ structure¹¹⁾⁻¹⁴⁾. Even if none of the known superstring theories turns out to be the TOE, some features of the above programme will probably survive. For example, we argue later that many forms of compactification⁴⁾ must lead to the same universal no-scale structure^{5),14)} at low energies. Crucial outstanding problems in any such scenario are the origin of supersymmetry breaking and determination of its scale, as well as the scales of compactification and of weak gauge symmetry breaking.

Many physicists conjecture that gaugino condensation¹⁵⁾ may trigger supersymmetry breaking, and it has recently been shown⁸⁾ how even if this does occur the cosmological constant may remain zero, at least in the absence of perturbative radiative corrections. Our main result in this paper is a characterization of the topological condition necessary for gaugino condensation arising from a twisted four-dimensional gauge field configuration in the compactified ten-dimensional field theory. The condition for bilinear gaugino condensation is non-trivial, and we give examples where known twisted gauge configurations only give higher order condensates. If gaugino condensation does occur in some hidden sector, and supersymmetry breaking is then fed through to the observable sector by higher order corrections, then the scales of supersymmetry and of gauge symmetry breaking can in principle be determined dynamically by radiative corrections, as proposed¹²⁾⁻¹⁴⁾ previously in the context of no-scale supersymmetric grand unified theories. Here the same radiative corrections would simultaneously determine dynamically the scales of compactification, of grand unification and of gaugino condensation:

$$\mu \sim \left(m_{3/2} / M_{Pl} \right)^{1/2} M_{Pl}, \quad m_X \sim \mu^4 / M_{Pl}^3, \quad \langle \bar{\chi}^i \chi^i \rangle \sim \mu^3 \quad (1)$$

No-scale models favour¹⁴⁾ sparticles weighing several hundred GeV, and the gravitino is likely to be even heavier, and hence to have sufficiently many decay modes $3/2 \rightarrow X\tilde{X}$ that it can decay before nucleosynthesis thus avoiding any cosmological gravitino problem^{16),17)}.

First we tackle the question of gaugino condensation, and determine circumstances under which it occurs. We work in the context of $E_8 \times E_8$ 'superstring theory'¹⁰⁾, but similar arguments could apply in a larger context. The mechanism proposed in Ref. 8) for the breakdown of supersymmetry with a vanishing cosmological constant involves a bilinear condensation of the gaugino field in E_8 . This condensation must occur in the full ten-dimensional theory in order for the mechanism to work. If the mechanism could be understood solely as an effect in the four-dimensional theory, the breaking of R-invariance in the four-dimensional Lagrangian would be spontaneous and not explicit. Then the term cM_{Pl}^3 in the superpotential, which is essential⁸⁾ for the vanishing of the cosmological constant, would be forbidden since it breaks the R-symmetry explicitly.

To obtain a ten-dimensional condensate $\langle \text{Tr}(\bar{\chi} \Gamma^{\alpha\beta\gamma} \chi) \rangle \neq 0$ we need topologically non-trivial gauge configurations (torons¹⁸⁾) which live in four dimensions. Torons are known to lead to $\langle \bar{q}q \rangle$ condensation in QCD¹⁹⁾ and to a $\langle \lambda\lambda \rangle$ condensate in four-dimensional supersymmetric Yang-Mills theory²⁰⁾, but the reasoning here is more subtle in several respects.

In four-dimensional SUSY Yang-Mills, the existence of a bilinear gluino condensate can already be inferred from the instanton computation²¹⁾. Instantons lead to $\langle \lambda\lambda(x)\lambda\lambda(y) \rangle \sim \Lambda^6$ where Λ is the scale parameter of the theory, and the quadrilinear condensate is independent of x and y , as a result of supersymmetry^{21),22)}. As x and y are infinitely separated, clustering can be used²²⁾ to conclude that we must have $\langle \lambda\lambda \rangle \sim \Lambda^3$, as confirmed by the toron computation²⁰⁾. In the ten-dimensional theory, the proliferation of indices makes it difficult to conclude on the basis of a ten-dimensional multilinear condensate that the bilinear condensate forms too. We shall see, however, that torons can generate directly the required $\langle \text{Tr}(\bar{\chi} \Gamma^{\alpha\beta\gamma} \chi) \rangle$.

We must allow for the possibility that the group E_8 is broken by the Hosotani mechanism²³⁾ to a subgroup Q , which is in general a direct product $Q = Q_1 \times Q_2 \times \dots \times Q_n$. Under E_8 , the gauginos transform as the adjoint representation which has dimension 248, and which happens to be also the fundamental representation of that group. Under the unbroken group Q , the 248 will transform as a direct sum of representations of the form $(r_1, \dots, r_n)$. We can only hope to achieve a condensation $\langle \text{Tr}(\bar{\chi}_i \Gamma^{\alpha\beta\gamma} \chi_i) \rangle \neq 0$ in a factor group Q_i for which the only massless mode in the direct sum is the adjoint representation, plus perhaps a bunch of singlets. If other representations r_i appear, the situation is rather

similar to massless four-dimensional SUSY QCD, where the existence of massless matter prevents the condensation of either $\lambda\lambda$ or $\bar{\psi}\psi$. This can be proved using supersymmetry and assuming no pathological behaviour of the condensates²²⁾, and is inevitable in any non-perturbative computation with topologically non-trivial gauge fields where both gauginos and matter fermions have zero modes: the resulting condensate cannot be clustered into condensates involving only one of those species. If, for example, E_8' is broken to $E_6 \times SU(3)$, such that both of the factor groups have massless matter in a fundamental representation, then no bilinear gaugino condensate forms. We therefore demand that there exists at least one Q_1 under which the gaugino χ is the only massless field transforming non-trivially. In particular, we retain for now the possibility that Q_1 could be the entire E_8' .

In the presence of a classical background gauge field A_μ^i in four dimensions, belonging only to the gauge group Q_1 , we wish to solve the ten-dimensional Dirac equation

$$\not{D}_{10} \chi = (\not{D}_6 + \not{D}_4) \chi = 0 \quad (2)$$

to find the zero modes χ_0^i . The four-dimensional covariant derivative contains the gauge connection A_μ^i and no spin connection, and the six-dimensional covariant derivative contains only a spin connection. One solution to Eq. (2) is readily found, namely

$$\chi_0^i = k \phi^{-3/4} F_{\mu\nu}^i \Gamma^{\mu\nu} \epsilon \quad (3)$$

where ϵ is the remaining $N = 1$ supersymmetry on the manifold $M \times K$, $\epsilon = \sum_{A=1}^4 \xi_A \times \eta_A$, where the ξ_A are constant spinors operating on four-dimensional space-time, and the η_A are covariantly constant spinors operating on the internal space. As was shown in Ref. 4), two independent η 's obtained with $SU(3)$ holonomy are in fact sufficient. To see that χ_0^i (3) is a zero mode, recall that if A_μ^i solves the classical equation of motion, so does its supersymmetric partner χ ²⁴⁾ so that $\not{D}_4 \chi_0^i = 0$. Then the only dependence of χ_0^i on the internal space co-ordinates comes from the η_A factors in ϵ and therefore also $\not{D}_6 \chi_0^i = 0$. The constant k is a normalization factor, which normalizes χ_0^i on ten-dimensional space-time. Another zero mode is trivially found from $\bar{\chi}_0^i$.

In the presence of a gauge configuration with arbitrary Pontryagin number, the solution (3) is not the only one which can be obtained for χ^1 . In order to get a bilinear condensate we have to make sure that there are no additional zero modes. This can be done when four-dimensional space-time M is compactified to a torus T^4 . Yang-Mills fields on the four-dimensional torus can have fractional Pontryagin number¹⁸⁾, and the Atiyah-Singer index theorem will then in some cases guarantee exactly two solutions to $D_4\chi = 0$. Any such solution could readily be made into a solution of $D_{10}\chi = 0$ using the covariantly constant spinors η_A as in Eq. (3).

Could the gauge group Q_1 be the full E_8 itself? The gauge group Q_1 , when represented in the adjoint representation of the group, forms a real vectorial fibre bundle on T^4 . The number of gaugino zero modes, according to the Atiyah-Singer theorem, will be given by the first Pontryagin number of this bundle. The Pontryagin number can be derived from the Pontryagin number of the associated principal fibre bundle. We find

$$P_{ad}/P_{prin} = C_2(ad) \quad (4)$$

where $C_2(ad)$ is the second Casimir of the adjoint representation, and $P_{ad} = 1/16\pi^2 \int \text{Tr}(F_{\mu\nu} \tilde{F}^{\mu\nu})_{ad}$. For $Q = E_8$, since $\Pi_1(E_8) = 0$ the classification of E_8 vector bundles on T^4 is the same as that of E_8 vector bundles on S^4 . The Pontryagin number of principal E_8 bundles on S^4 is given by the winding number $\Pi_3(E_8) = \mathbb{Z}^*$, which implies that the minimal number of gaugino zero modes is $C_2(ad) = 60$. The resulting multiple gaugino condensate would contain 60 gaugino fields, which is not exactly what we want and we therefore conclude that E_8 has to be broken in order to obtain two zero modes.

Could the gauge group Q_1 be an $SU(N)$ subgroup of E_8 ? When $SU(N)$ is represented in the adjoint representation, the resulting vector bundle is a $PU(N) = SU(N)/Z(N)$ bundle. Since $\Pi_3(PU(N)) = \mathbb{Z}$ and since $C_2(ad) = 2N$, the regular instanton bundle will have $2N$ gaugino modes. However, $\Pi_1(PU(N)) = \mathbb{Z}(N)$, and therefore the classification of $PU(N)$ bundles on T^4 is different from that on S^4 . In fact, we can get non-trivial classical gauge configurations, the torons¹⁸⁾, with $P_{ad} = 2$ (or equivalently $P_{fund} = 1/N$). These configurations will

*) Note that $\Pi_9(E_8) = 0$ which means that E_8 has no winding number on S^{10} , and therefore no non-perturbative θ vacuum angle in the theory before compactification.

lead to a bilinear gaugino condensate which has a finite limit as the volume of the torus is taken to infinity. The full computation of the condensate is given in Ref. 20).

In order to use the toron mechanism, we have seen that we have to break E_8' to a group containing an $SU(N)$ factor, say

$$E_8' \rightarrow SU(5) \times SU(5), SU(3) \times E_6, SU(2) \times E_7, \text{ etc.} \quad (5)$$

Under all the possible breakings of E_8' , the 248 representation splits into representations of the Q_i which always include at least one non-trivial matter representation for each of the Q_i products. As explained above, these representations have to be disposed of, if a bilinear gaugino condensate is to form. A zero mode of $\mathcal{V}_4$ will be a zero mode of $\mathcal{V}_{10}$ only if it is also a zero mode of $\mathcal{V}_6$, namely if it is a zero mode of the internal space. These transform in a characteristic way under the discrete symmetry group G , which determines the internal space manifold K/G . On any manifold with $SU(3)$ holonomy, there is one solution of the type $(0,0)$ and one of the type $(0,3)^{6)*}$.

The Hosotani mechanism²³⁾ maps the elements $y \in G$ into $U_g' \in E_8'$ such that $g \rightarrow U_g'$ is a homomorphism of G into $E_8'^{6)}$. The unbroken subgroup of E_8' will be that which commutes with all the U_g' . The U_g' themselves form a discrete subgroup of E_8' , which is denoted by $\bar{G}'$. Arguing on the basis of Witten's analysis⁶⁾, the only representations of E_8' which survive as zero modes on the manifold K/G are those which are singlets under the diagonal sum $G \oplus \bar{G}'$. This enables us to dispose of unwanted zero modes.

Take for example the manifold⁴⁾ $K = Y_{(4;5)}$, the solution of a quintic equation in CP^4 . The discrete symmetry of the manifold is $G = Z(5) \times Z(5)$, and the forms $(0,0)$ and $(0,3)$ are both invariant under G . This is the general condition for retaining $N = 1$ supersymmetry on K/G . Let us now choose $\bar{G}'$ to be also $Z(5) \times Z(5)$. This will induce the breaking $E_8' \rightarrow SU(5) \times SU(5)$ under which the adjoint of E_8' has the decomposition

$$\underline{248} = (\underline{24}, \underline{1}) + (\underline{1}, \underline{24}) + (\underline{10}, \underline{5}) + (\underline{10}, \underline{5}) + (\underline{5}, \underline{10}) + (\underline{5}, \underline{10}) \quad (6)$$

*) A differential form of the type (p,q) has p holomorphic and q antiholomorphic indices.

Since the solutions of $D_6 \chi = 0$ on K are invariant under G , they must be invariant under $\bar{G}'$ in order to remain solutions on K/G . Obviously, the only surviving representations are $(1,24)$ and $(24,1)$. The forms $(0,0)$ and $(0,3)$ for these representations will combine to give one Majorana-Weyl spinor in ten dimensions, the ε of Eq. (3). When we introduce a toron in either of the $SU(5)$ factors, this will lead to two gaugino zero modes. The gauge invariant condensate in the ten-dimensional theory is $\langle \text{Tr}(\bar{\chi} \Gamma_{\alpha\beta\gamma} \chi) \rangle$, which we now compute.

The ten-dimensional field strength $F_{\mu\nu}^{(10)}$ to be used in Eq. (3) is obtained from the four-dimensional toron $F_{\mu\nu}^{(4)}$ via the Weyl rescaling of the metric²⁵⁾

$$F_{\mu\nu}^{(10)} = e^{3\sigma} F_{\mu\nu}^{(4)} = e^{3\sigma} F_{\mu\nu}^{\text{toron}} \quad (7)$$

The action of the toron is

$$\begin{aligned} S^{\text{toron}} &= \int \frac{\phi^{-3/4}}{4g_{10}^2} \text{Tr} F_{\mu\nu}^{(10)} F^{\mu\nu(10)} \sqrt{\det g^{(10)}} d^{10}x \\ &= \frac{1}{4} \int \phi^{-3/4} e^{3\sigma} F_{\mu\nu}^{(4)} F^{\mu\nu(4)} d^4x \end{aligned} \quad (8)$$

with the choice $g_{10}^2 M_{\text{Pl}}^6 = 1$. For an $SU(N)$ gauge theory,

$$S^{\text{toron}} = 8\pi^2/g_4^2 N \quad (9)$$

The toron has four translation gauge zero modes A_0^{toron} , whose Jacobian contributes to the condensate a factor [see Ref. 26)]

$$\begin{aligned} &\frac{d}{g_4^2} \mu^4 \left[\left\{ \int (A_0^{(10)})^2 \sqrt{\det g^{(10)}} d^{10}x \right\}^{1/2} \right]^4 \\ &= d \cdot \mu^4 \left[\int d^4x (A_0^{\text{toron}})^2 \right]^2 \end{aligned} \quad (10)$$

where we have again used the Weyl rescaling²⁵⁾, d is a constant of order 1, and μ is the scale at which the four-dimensional coupling constant is $g_4^2 = \phi^{3/4} e^{-3\sigma}$, namely the compactification scale. The two Majorana-Weyl zero modes contribute an additional factor μ^{-1} , and the determinant of non-zero modes is 1 because of supersymmetry. We find:

$$\begin{aligned} & \langle \text{Tr } \bar{\chi}^i \Gamma_{\alpha\beta\gamma} \chi^i \rangle \\ & \sim \mu^3 \left[\int d^4x (A_0^{\text{toron}})^2 \right]^2 \int d^4z \text{Tr} \langle \bar{\chi}_0^i \Gamma_{\alpha\beta\gamma} \chi_0^i(x-z) \rangle \times e^{-S_{\text{toron}}} \end{aligned} \quad (11)$$

where z is the "centre" of the toron. The four-dimensional theory sees the piece $\langle \text{Tr}(\bar{\chi} \Gamma_{ijk} \chi) \rangle = A \epsilon_{ijk}$. Note that even though i, j and k are indices on the internal manifold, such a condensate can be realized in terms of the four-dimensional solution χ_0^i of Eq. (3).

Now that we have a gaugino condensate, we indicate how the scales of compactification, of supersymmetry breaking and of weak gauge symmetry breaking may be dynamically determined by radiative corrections. As was shown in Refs. 5) and 8), the effective four-dimensional theory at energies below the compactification scale is an $N = 1$ supergravity theory containing two gauge singlet chiral superfields

$$S \equiv e^{3\sigma} \phi^{-3/4} - i \frac{2b_0}{3} \alpha \equiv S_R + i S_I \quad (12)$$

$$T \equiv e^{\sigma} \phi^{3/4} + i a \equiv T_R + i T_I \quad (13)$$

where ϕ is the dilaton in the supergravity multiplet, σ determines the scale of the six compactified dimensions:

$$g_{\text{IS}} = e^{\sigma} g_{\text{IS}}^0 : \int_K d^6x \sqrt{\det g^0} = M_{Pl}^{-6} \quad (14)$$

α is the phase of the gaugino condensate, and a is a component of the second rank tensor field, as well as gauge non-singlet chiral superfields denoted generically by ϕ^i . Couplings of the chiral superfields are characterized⁵⁾ by a Kähler potential

$$G = -\ln(S+S^\dagger) - 3\ln(T+T^\dagger - \phi^i \phi_i^\dagger) + \ln |W|^2 \quad (15)$$

where $W(S, \phi^i)$ is the chiral superpotential, and by a particularly simple form⁸⁾ of the chiral function characterizing gauge kinetic terms:

$$\int d^2\theta f_{\alpha\beta} W^\alpha W^\beta + \text{h.c.} \quad , \quad f_{\alpha\beta} = \delta_{\alpha\beta} S \quad (16)$$

As was shown in Ref. 8), condensation in the hidden sector makes the following contribution to the superpotential:

$$W(S) = M_{Pl}^3 \left(c + h \exp \frac{-3S}{2b_0} \right) \quad (17)$$

in the notation of Ref. 8). Because of the "no-scale" $SU(n,1)/SU(n) \times U(1)$ form¹⁴⁾ of the T - and ϕ^1 -dependent terms in G (15), to which we return later, the S superpotential (17) yields a positive semi-definite effective potential

$$V = \frac{1}{16} M_{Pl}^4 \frac{1}{S_R T_R^3} \left| c + h \left(\frac{3S_R}{T_R} + 1 \right) \exp \frac{-3S}{2b_0} \right|^2 \quad (18)$$

As long as $c \neq 0$, the potential (18) may be minimized at $V = 0$ by a suitable dynamical determination of S . The phase α is chosen so that the c and h terms in (18) have opposite signs, and then $S_R = e^{3\sigma} \phi^{-3/4}$ is fixed so that $S_R = O(1)$. The gravitino then acquires a mass

$$m_{3/2} = M_{Pl} e^{G/2} \sim M_{Pl}^3 \frac{|c|}{S_R^{1/2} T_R^{3/2}} = O(M_{Pl}) \times e^{-6\sigma} \times |c| \quad (19)$$

and the mass of the axion associated with the phase α is comparable. Since h which is computed from the gaugino condensate is of order 1, c will typically be a small number: $c = O(1) \exp - 2\pi^2/g^2 N$. Note, however, that so far T_R and hence $e^\sigma = (S_R T_R)^{1/4}$ are undetermined.

To determine them we turn to the observable sector ϕ^1 . The origin of their "no-scale" structure¹⁴⁾ in the Kähler potential (15) is easy to see and likely to be a general feature of models compactified from higher dimensions. A general feature of such models is that one wants the four-dimensional cosmological constant Λ to be (almost) zero, and there to exist a considerable number of massless ($m \ll M_{Pl}$) chiral superfields. These parametrize the continuous degrees of freedom characterizing a family of related compactification manifolds, each of which has $\Lambda = 0$ by construction. Since these models have the same vacuum energy, the values of the chiral superfields are undetermined at this level and remain to be determined by radiative corrections. The only known Kähler potentials which give $\Lambda = 0$ automatically for such a continuous range of field

variables are those containing a $-3\ln(T+T^+-\phi^1_1\phi^+_1)$, i.e. an $SU(n,1)/SU(n)\times U(1)$ structure¹⁴⁾, in the context of which just such a scenario for the dynamical determination of light mass scales has been proposed¹²⁾⁻¹⁴⁾.

Here, supersymmetry breaking originates from gaugino condensation in the hidden sector, and is absent from the observable sector at tree level, where it can only arise from field-theoretical loops cut-off at the scale of the condensate. It therefore seems quite possible that the observable supersymmetry breaking mass parameters m_0 (for sleptons) and $m_{1/2}$ (for gauginos) may be somewhat less than $m_{3/2}$, though this requires further investigation. Exploiting such a seed of supersymmetry breaking, weak gauge symmetry breaking can be triggered in no-scale models¹⁴⁾ by a heavy top or heavier quark²⁷⁾, whose large Yukawa coupling renormalizes the Higgs mass squared m_H^2 so that although positive (or zero) at the scale of the condensate where m_0 and $m_{1/2}$ are generated, it finally becomes negative at some renormalization scale

$$\mu_0 = \mu \exp -\frac{O(1)}{\alpha} \quad (20)$$

which can easily be $O(100)$ GeV. The scale μ_0 is independent of $m_{0,1/2}$ as long as $\mu_0 \gg m_{0,1/2}$. As soon as $m_H^2 < 0$ at scales $\sim \mu_0$, weak gauge symmetry breaking is possible with a scale determined by dimensional transmutation²⁸⁾:

$$m_w \propto \langle 0|H|0 \rangle = O(\mu_0) \quad (21)$$

In models such as the present one where the scales $m_{0,1/2}$ of global supersymmetry breaking are as yet undetermined, they too can be determined dynamically¹²⁾:

$$m_{0,1/2} = O(\mu_0) \propto m_w \quad (22)$$

The steps in Eqs. (20) to (22) are familiar to experts in the field of supergravity model-building: next comes a novelty. By virtue of (22), we now determine through Eq. (19):

$$T_R = O(m_{3/2}^{-2/3}) \quad (M_P \equiv 1) \quad (23)$$

and hence

$$e^{\sigma} = O(m_{3/2}^{-1/6}) \quad (24)$$

Thus the scale of compactification is also dynamically determined:

$$\mu = e^{-\sigma/2} = O(m_{3/2}^{1/12}) \quad (25)$$

This should come as no surprise, since by having a family of compactification manifolds with $\Lambda = 0$ at the tree level one leaves the scale of compactification free to be determined by radiative corrections. Much has been made in the past of the fact²⁹⁾ that perhaps $m_X \ll M_{Pl}$. Here we find that

$$m_X \sim \mu^4 / M_{Pl}^3 \sim e^{-2\sigma} \sim O(m_{3/2}^{1/3}) \quad (26)$$

and is also dynamically determined. This is because, although the scale of the Hosotani symmetry breaking mechanism is μ , the gauge boson masses that it provides^{8), 25)} are $m_X = O(\mu^4 / M_{Pl}^3)$. Above this scale, the couplings of all the gauge bosons evolve in an identical unified way.

Finally, we make a cosmological comment. One of the cosmological problems which have bedevilled supersymmetric model-builders has been the gravitino problem¹⁶⁾: how to avoid gravitino decays from spoiling the successes of conventional primordial nucleosynthesis. Here it is natural to suppose that the gravitino is massive enough to decay before nucleosynthesis, thereby avoiding the problem. We have the suggestion that $m_{3/2} \gg m_{0,1/2}$, and very possibly $m_{3/2} > m_{\tilde{q}, \tilde{\ell}, \tilde{W}, \tilde{B}, \tilde{g}}$. The gravitino therefore has many decay modes open to it, and its lifetime is

$$\tau_{3/2} \approx \frac{4 \times 10^8 \text{ sec}}{12 + 2N_{\phi}} \times \left(\frac{100 \text{ GeV}}{m_{3/2}} \right)^3 \quad (27)$$

where N_{ϕ} is the number of light chiral superfields: $N_{\phi} = 49$ in the minimal standard model. The gravitino will indeed decay sufficiently rapidly ($\tau_{3/2} < 100\text{s}$) if

$$\left(m_{3/2}/100\text{GeV}\right) > \frac{74}{(12+2N_\phi)^{1/3}} \quad (28)$$

which is very possibly the case. We conclude that the gravitino problem is potentially solved.

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