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GUTS AND SUPERSYMMETRIC GUTS IN THE VERY EARLY UNIVERSE*

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0 INTRODUCTION

This talk is intended as background material for many of the other talks treating the possible applications of GUTs to the very early universe. I start with a review of the present theoretical and phenomenological status of GUTs before going on to raise some new issues for their prospective cosmological applications which arise in supersymmetric (susy) GUTs. The first section is an update on conventional GUTs (Ellis 1981), which is followed by a reminder of some of the motivations for going supersymmetric (Dimopoulos & Georgi 1981; Sakai 1982). There then follows a simple primer on susy and a discussion of the structure and phenomenology of simple susy GUTs. Finally we come to the cosmological issues, including problems arising from the degeneracy of susy minima, baryosynthesis and supersymmetric inflation, the possibility that gravity is an essential complication in constructing susy GUTs and discussing their cosmology, and the related question of what mass range is allowed for the gravitino. Several parts of this write-up contain new material which has emerged either during the Workshop or subsequently. They are included here for completeness and the convenience of the prospective reader. Wherever possible, these anachronisms will be flagged so as to keep straight the historical record.

1 STANDARD GUTS

Presumably you are familiar with the motivations and guiding principles of conventional GUTs (Georgi & Glashow 1974; Georgi et al. 1974). They will not be discussed here, but just an update given on their phenomenological status (for more details see Ellis 1981). GUTs make several low energy predictions, some of which work and some of which are less successful. First the good news: standard GUTs predict (Marciano & Sirlin 1981; Llewellyn Smith et al. 1981)

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$$\sin^2\theta_W^{eff}(\langle Q^2\rangle = 20\ \mathrm{GeV}^2) = 0.215 \pm 0.002 \qquad (1)$$

for $\Lambda_{\overline{MS}}$ (4 operational flavors) = 0.1 to 0.2 GeV. This is to be compared with the experimental value for the effective value of $\sin^2\theta_W$ (including radiative corrections) at an average momentum transfer $\langle Q^2\rangle$ = 20 GeV^2 in deep inelastic νN scattering (Marciano & Sirlin 1980; Llewellyn Smith & Wheater 1981):

$$\sin^2\theta_W^{eff}(\langle Q^2\rangle = 20\ \mathrm{GeV}^2) = 0.216 \pm 0.012 \quad . \qquad (2)$$

Another successful prediction (Chanowitz <u>et al</u>. 1977; Buras <u>et al</u>. 1978; Nanopoulos & Ross 1982) is for the b quark mass, deduced from the τ lepton mass:

$$m_b \approx 5\ \mathrm{GeV} \qquad (3)$$

if one assumes the existence of 6 quark flavors corresponding to 3 light neutrinos. Now for the <u>less good news</u>: in the same way that (3) was deduced from m_τ one can also deduce from m_μ that (Buras et al 1978):

$$m_s \approx \tfrac{1}{2}\ \mathrm{GeV} \quad (?) \qquad (4)$$

whereas many theorists believe that the true short distance value of $m_s \approx 150$ MeV (Weinberg 1977). I am not convinced by their arguments and prefer to wait and see what value of m_s will emerge from lattice QCD calculations. Finally the <u>bad news</u>: on the same basis as the predictions (3) and (4) one can also predict (Buras <u>et al</u>. 1978)

$$m_d/m_s = m_e/m_\mu \qquad (5)$$

while experimentally we know that $m_e/m_\mu \approx 1/200$ and conventional QCD phenomenology suggests that $m_d/m_s \approx 1/20$. Despite the quantitative difficulties (4,5) the quark and lepton masses are qualitatively in the correct relationship, and perhaps some small effect can come in at the level of MeV to cure the ratio (5) for m_d/m_e (Ellis & Gaillard 1979).

Thus, undiscouraged, we do not yet loss faith in the most existing prediction of conventional GUTs, namely for baryon decay. From Figure 1 we find a decay amplitude $\propto g^2/m_X^2$, where g is the GUT coupling constant and m_X the superheavy boson mass, from which we deduce

$$\Gamma(B \to \bar{\ell} + X) \propto 1/m_X^4 \quad , \quad \tau_B \propto m_X^4 \quad . \qquad (6)$$

In simple models such as minimal SU(5) or SO(10) one can estimate

$$m_X \simeq (1\ \mathrm{to}\ 2) \times 10^{15} \times \Lambda_{\overline{MS}} \qquad (7)$$

where $\Lambda_{\overline{MS}}$ is the strong interaction scale parameter estimated to be

$$\Lambda_{\overline{MS}} \simeq (0.1 \text{ to } 0.2) \text{ GeV} \quad . \tag{8}$$

Combining Equations (7) and (8) we infer that

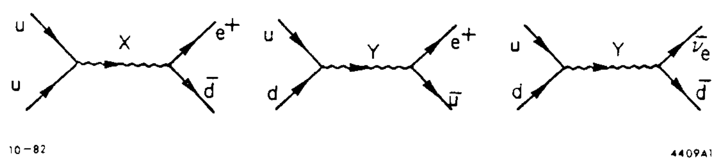
$$m_X = (1 \text{ to } 4) \times 10^{14} \text{ GeV} \quad . \tag{9}$$

Various estimates of the constants of proportionality in the dependence (6) on m_X then suggest (Ellis et al. 1980) that

$$\tau_B \simeq (10^{27} \text{ to } 10^{31}) \text{ years} \tag{10}$$

in simple GUTS. There are now two experiments (Krishnaswamy et al. 1981, 1982; Battistoni et al. 1982) reporting positive evidence for baryon decay — one since the Workshop-corresponding to a nucleon lifetime of a few times 10^{30} years. It will be interesting to see whether these preliminary indications are confirmed, and whether the branching ratios for baryon decay conform to the conventional GUT predictions. If so, it would be a dramatic confirmation of GUT ideas.

Fig. 1. Diagrams contributing to the dimension 6 operators responsible for baryon decay in minimal SU(5).



Ironically, even as our experimental colleagues make these fascinating observations, theoretical fashion is deserting conventional GUTs. The reason is the hierarchy problem: why/how is $m_W/m_X << 1$? Conventional GUTs contain two vastly different energy scales m_W and m_X, and m_W in particular is much less than the natural candidate for a basic mass scale, namely the Planck mass $m_P \simeq 10^{19}$ GeV. In conventional GUTs the two mass scales m_W and m_X are associated with two sets of Higgs fields -- e.g. the $\underline{24}$ ϕ and the $\underline{5}$ H of Higgses in minimal SU(5). These are to have a ratio of vacuum expectation values of order 10^{-12}:

$$m_W = \frac{gv}{2} \, , \; m_X = \frac{5gV}{2\sqrt{2}} \; : \; \langle 0|H|0\rangle \equiv v \, , \; \langle 0|\phi|0\rangle \equiv \frac{V}{\sqrt{30}} \begin{pmatrix} 2 & & & & 0 \\ & 2 & & & \\ & & 2 & & \\ & & & -3 & \\ 0 & & & & -3 \end{pmatrix} \tag{11}$$

To achieve this gymnastic feat we need a scalar potential $V(\phi,H)$ in which the effective mass of the SU(2) doublet components of the H multiplet is very small:

$$|m_H^2| \lesssim O(10^{-24})\, V^2 \quad . \tag{12}$$

This feat is difficult to arrange, since there are many contributions to the scalar masses which are a priori much larger than (12), suggesting the necessity of some concealed symmetry if bizarre and unmotivated cancellations are to be avoided. For example, there is the effect of propagating scalar particles through space-time foam (see Figure 2(a)):

$$\delta m_H^2 = O(m_P^2) \tag{13}$$

according to the local experts (Hawking et al. 1979, 1980). Then there are quadratic divergences in perturbation theory (see Figure 2(b)) which yield

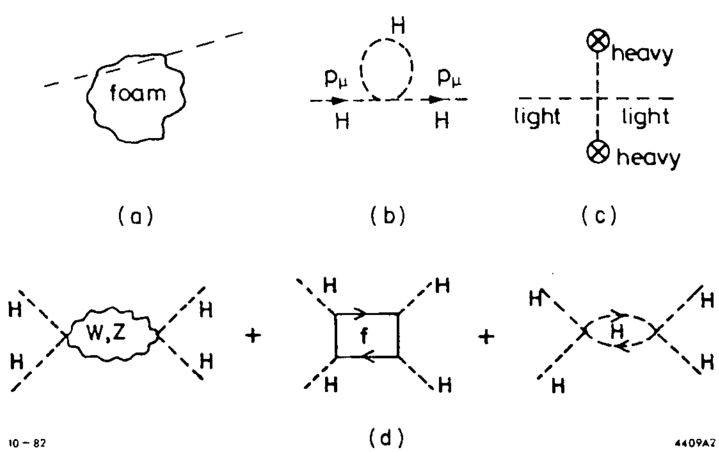
$$\delta m_H^2 \propto \int^{\Lambda} d^4k \, \frac{1}{k^2} \propto \Lambda^2 : \Lambda^2 = O\left(m_X^2 , m_P^2\right) ? \tag{14}$$

There are also contributions to the light Higgs masses from the Higgses with large vacuum expectation values (see Figure 2(c)):

$$\delta m_H^2 \propto V^2 \quad . \tag{15}$$

Interactions like the $\phi^2 H^2$ shown in Figure 2(c) are generated by radiative corrections (see Figure 2(d)) even if they were unlawfully excluded from the tree-level Lagrangian (Gildener 1976, Buras et al. 1978). One must adjust the parameters of the Lagrangian taking into account at least 12

Fig. 2. (a) A scalar particle propagating through space-time foam may acquire a mass $O(m_P)$ (Hawking et al. 1979, 1980). (b) A quadratically divergent contribution to the scalar boson (mass)^2. (c) A large contribution to the "light" Higgs (mass)^2 from the condensation of "heavy" Higgs in the vacuum. (d) One-loop contributions to the Higgs self-couplings.



loop diagrams in perturbation theory so as to reduce the right-hand-side of Equation (15) by $O(\alpha^{12}) = O(10^{-24})$ to make it acceptably small.

Theorists have toyed with several possible solutions to this hierarchy problem. The latest and most promising is supersymmetry (susy), so let us now look at the construction of susy GUTs.

2 SUPERSYMMETRY PRIMER

Before constructing susy GUTs, let us first refresh our recollections of susy and of its phenomenology. What is susy? It is a symmetry between fermions and bosons which is generated by charges Q with spin ½ (Gol'fand & Likhtman 1971; Volkov & Akulov 1973; Wess & Zumino 1974a). They obey an algebra of anticommutators, as one would expect for fermion operators, namely

$$\left\{Q^i_\alpha , Q^{\dagger\beta}_j\right\} = (\sigma^\mu)^\beta_\alpha \, P_\mu \, \delta^i_j \tag{16}$$

where the $i = 1,\ldots,N$ are the different types of extended supersymmetry, the α and β are spinorial indices, and P_μ is the momentum operator. Each operator Q^i_α changes spin by half a unit. Since gauge theories only have spins (or helicities) between ±1, they can accommodate at most $N = 4$ (global) supersymmetries. If one makes the supersymmetry transformations local, then Equation (16) means that one necessarily must include general coordinate transformations. Therefore one must include gravity and accommodate spins between ±2, in which case $N \leq 8$ supersymmetries are allowed (see Section 6) (Van Nieuwenhuizen 1981). In all these cases one has equal numbers of boson and fermion states $|B\rangle$ and $|F\rangle$, since

$$Q|B\rangle = |F\rangle \quad \text{and} \quad Q|F\rangle = |B\rangle \tag{17}$$

respectively. We will restrict ourselves to simple $N = 1$ susy in the discussions that follow. In this case there are two classes of supermultiplets (representations of the global supersymmetry algebra) which are relevant. They are the

$$\text{gauge supermultiplet} : \begin{pmatrix} 1 \\ \tfrac{1}{2} \end{pmatrix} \tag{18}$$

which must lie in an adjoint representation of the gauge group, and the

$$\text{chiral supermultiplet} : \begin{pmatrix} \tfrac{1}{2} \\ 0,0 \end{pmatrix} \tag{19}$$

which may lie in any representation of the gauge group. If $N \geq 2$ we are forced to have fermions lying in real representations of the gauge group, which conflicts with the phenomenological (and GUT) necessity of a complex fermion representation.

Why susy? This is a question with many answers, one of which is simply because susy is beautiful. But this is a rather subjective argument -- mother aardvarks also believe that their offspring are beautiful, a point of view that can be debated. Susy is the only known type of symmetry which has not yet been exploited in fundamental physics, so perhaps we should look for an application. It reduces the number of divergences in quantum field theory: for example the $N = 4$ gauge theory appears to be completely finite (Mandelstam 1982), while the $N = 8$ supergravity theory is probably finite up to 6 loops (Grisaru & Siegel 1982). Susy provides a home for scalar fields, as can be seen from the chiral supermultiplet (19) where they are linked together with spin ½ fermions. Susy has the general property of linking together matter and radiation -- witness the gauge supermultiplet (18) relating spin 1 gauge bosons to spin ½ fermions, and the $N = 1$ supermultiplet containing the graviton and the gravitino discussed in Section 6. Extended supergravity theories are the only available candidates for unifying particle physics and gravity, though it must be confessed that present models still leave something to be desired (Ellis et al. 1982b).

All the previous motivations are somewhat philosophical. The practical reason for much current interest in susy is the prospect of "solving" the hierarchy problem in GUTs (Dimopoulos & Georgi 1981; Sakai 1982). "Solving" is in quotation makes because susy does not (yet) cast any convincing light on the question why $m_W/m_P \simeq O(10^{-17})$ or $m_W/m_X \simeq O(10^{-13})$, but it can alleviate the technical difficulties of maintaining this hierarchy once it has been imposed. This is because of the reduction in the number of divergences mentioned earlier. There are no quadratic divergences in susy theories, which removes one of the worst corrections (14) to scalar (Higgs) boson masses. Also, the only logarithmic divergences are wave function and gauge coupling renormalizations. There are no intrinsic renormalizations of the Yukawa couplings or the Higgs self-couplings (Wess & Zumino 1974b; Iliopoulos & Zumino 1974; Ferrara et al. 1974). This means that if one sets the contribution (15) to m_H^2 equal to zero at the tree level, it will not be generated by radiative corrections.

Perhaps we will eventually need to worry about quantum gravity (13) as well, but at least some progress has been made.

How susy? Susy cannot be exact in the real world, as it would imply equal masses for bosons and fermions $m_B = m_F$, just as conventional isospin invariance implies $m_p = m_n$. One must therefore ask how badly susy can be broken, i.e., how heavy can unseen supersymmetric partners be? Presumably the scale of susy breaking must be less than the scale $O(m_P \approx 10^{19}$ GeV) associated with gravitation. A more stringent constraint is that of "solving" the hierarchy problem. Let us suppose God in her infinite wisdom fixes $m_{Higgs} = 0$ (why? this is the reason that the hierarchy problem is not really solved). In an exactly susy world m_{Higgs} would stay zero, but if susy is broken there are corrections due to imperfect cancellations between the diagrams shown in Figure 3:

$$\delta m_H^2 \simeq O(1/16\pi^2)(g^2 \text{ or } \lambda^2)(m_B^2 - m_F^2) \quad (20)$$

where $g(\lambda)$ is a gauge (Yukawa) coupling. In order to maintain the Weinberg-Salam Higgs boson sufficiently light: $|\delta m_H^2| \lesssim O(1)$ TeV^2 we must require

$$|m_B^2 - m_F^2| \lesssim \frac{O(16\pi^2)}{(g^2 \text{ or } \lambda^2)} \times O(1) \text{ TeV}^2 \quad . \quad (21)$$

Thus we see that while the susy breaking may be relatively small, it can be considerably larger than O(1) TeV^2 in a supermultiplet which is weakly coupled to the Weinberg-Salam Higgs multiplet ($\lambda \ll 1$) (Ellis et al. 1982c). There is no necessity for the susy breaking mass splittings $|m_B^2 - m_F^2|$ to be universal for different supermultiplets.

In all theories of broken susy we expect to find the light susy partner particles listed in Table I. Alongside each sparticle entry is indicated the most stringent present lower limit on its mass. All of

Fig. 3. One-loop contributions to the (mass)^2 of scalar particles.

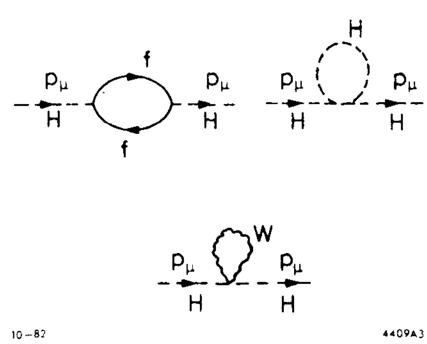


Table I. Supersymmetric Particles

	Particle	Partner	Spin	Lower Limit on Mass
Gauge Supermultiplets	gluon g	gluino $\tilde{g}$	½	3 GeV
	$W^{\pm}$	wino $\tilde{W}^{\pm}$	½	15 GeV
	γ, Z	photino $\tilde{\gamma}$, zino $\tilde{Z}$ or wino $\tilde{W}^{o}$, bino $\tilde{B}$	½	0 GeV
Chiral Supermultiplets	quarks q	squarks $\tilde{q}$	0	15 GeV
	leptons ℓ	sleptons $\tilde{\ell}$	0	15 GeV
	Higgses H	shiggses $\tilde{H}^{\pm}$	½	15 GeV
		$\tilde{H}^{o}$	½	0 GeV

the lower limits come from unsuccessful searches for new particles at PETRA, with the exception of the gluino mass limit which comes from hadron-hadron collisions and perhaps heavy quarkonium decays (Farrar & Fayet 1978 a,b; Campbell et al. 1982).

There are several possible philosophies for breaking susy. One is to do it explicitly, introducing arbitrary "soft" mass terms for unseen particles which are sufficient to push them above the lower bounds in Table I. This strategy is aesthetically unattractive. It introduces many new parameters whose magnetude is not explained, and does not respect naturally the constraints imposed by flavor-changing neutral interactions (Ellis & Nanopoulos 1982a) and by CP violation (Ellis et al. 1982a). It is more appealing to break susy spontaneously.

Either of two strategies may be followed. One approach has spontaneous susy breaking arising in the gauge sector, and is called D-breaking (Fayet & Iliopoulos 1975). In the other approach soft susy breaking arises in the Yukawa interactions of the chiral superfields (F-breaking) (O'Raifeartaigh 1975; Fayet 1975). These names derive from the conventional notation for the scalar field potential (Fayet & Ferrara 1977)

$$V = \sum_{\alpha} |D_{\alpha}|^2 + \sum_{i} |F_{\phi_i}|^2 \tag{22}$$

where

$$D_{\alpha} \equiv g_{\alpha} \sum_{i} \bar{\phi}_i T_{\alpha} \phi_i + \xi_{\alpha} \tag{23}$$

with the g_α the different gauge couplings and ξ_α an arbitrary constant appearing if α is a U(1) factor, and

$$F_{\phi_i} \equiv \partial P/\partial\phi_i \tag{24}$$

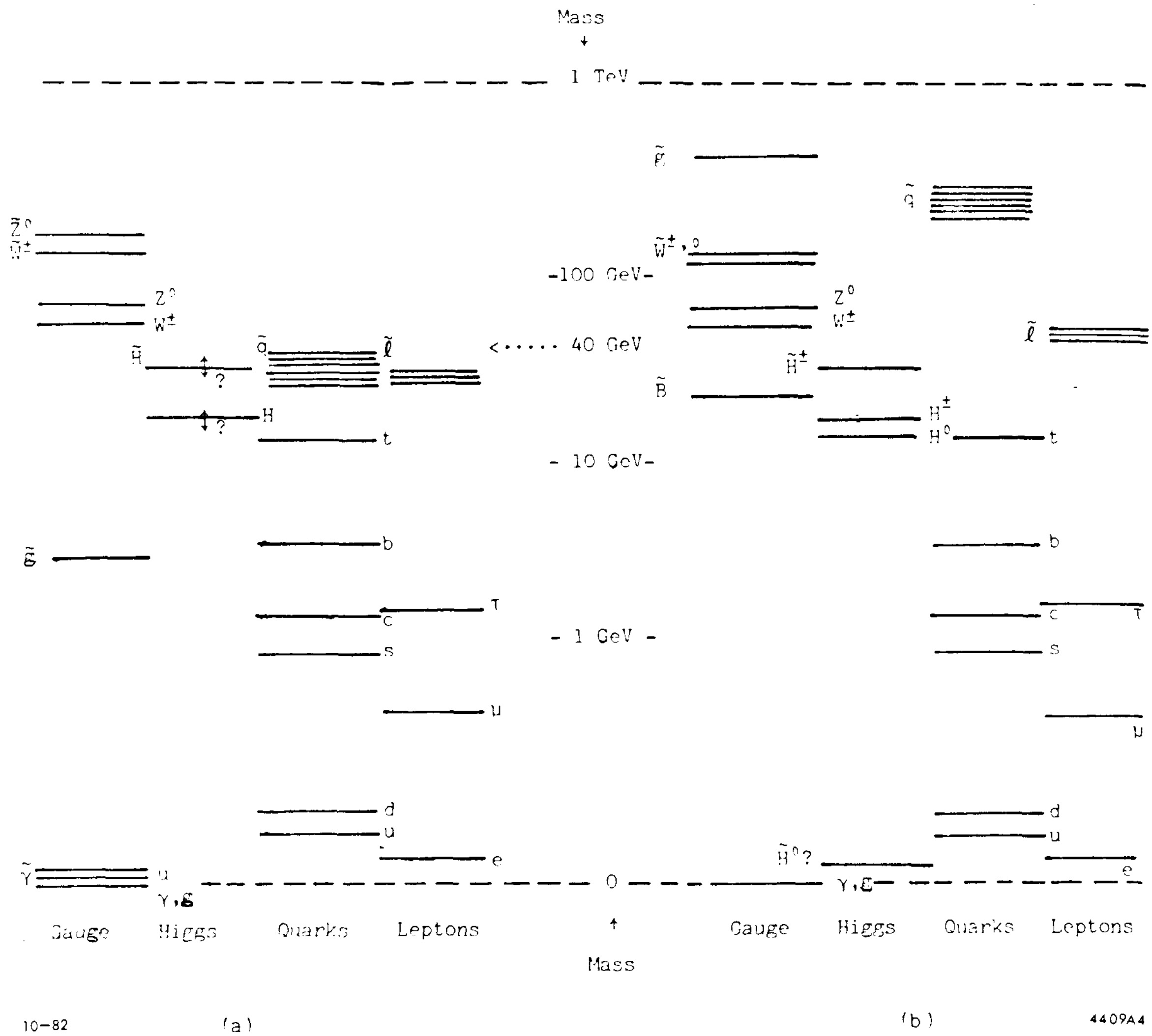
where

$$P = \sum_{i,j,k} \lambda_{ijk}\phi_i\phi_j\phi_k + \sum_{i,j} m_{ij}\phi_i\phi_j + \sum_i a_i\phi_i + b \tag{25}$$

is a cubic polynomial in the chiral superfields ϕ_i called the superpotential. D-breaking models require a new U(1) factor in the gauge group with an associated gauge boson U, and attempt to get realistic mass spectra at the tree level (Fayet 1981). F-breaking models rely on radiative corrections to feed the susy breaking through to all the known supermultiplets of Table I, and are technically more complicated to analyze.

These two classes of models have rather different phenomenologies, and it is advisable to keep both in mind. In the D-breaking models one can derive some upper limits of order 40 GeV on the masses of squarks and sleptons (Fayet 1981), but only under simplifying assumptions which need not be valid. Conventional neutral current phenomenology forbids $m_U/m_{Z^0} = O(1)$, but $m_U \ll m_{Z^0}$ and $m_U \gg m_{Z^0}$ are both tenable hypotheses (Fayet 1981; Barbieri et al. 1982 a,b). It is a noteworthy feature of F-breaking models that the sparticles of Table I are naturally heavier than their familiar partners. Gauginos are heavier than gauge bosons because they are not protected by gauge invariance, while squarks and sleptons are heavier than quarks and leptons because they are not protected by chiral symmetry. In either class of models the primordial scale of m_S susy breaking could in principle be as large as $O(m_X$ or $m_P)$, though we will see later that this may give problems when we make susy local and go to a supergravity theory. Figure 4(a),(b) illustrate two possible scenarios for the spectroscopy of "light" susy particles, based on D- and F-breaking models respectively (Fayet 1981; Ellis et al. 1982 c,d).

Fig. 4. Possible spectroscopies (a) in susy broken à la Fayet & Iliopoulos (1975), and (b) in susy broken à la Fayet (1975) and O'Raifeartaigh (1975).



3 SUPERSYMMETRIC GUTS

Now that we have developed some notions about the spectroscopy of spontaneously broken susy, we will now proceed to construct simple susy GUTs and discuss their phenomenology. The fermions and Higgs of conventional GUTs are now assigned to chiral supermultiplets. For example, the conventional $\underline{\overline{5}}$ ψ^{α} and $\underline{10}$ $\chi_{\alpha\beta}$ of fermions in a normal SU(5) generation now become chiral supermultiplets with additional spin-zero components (see Equation (19)). Likewise the adjoint $\underline{24}$ of Higgses ϕ^{α}_{β} also acquires spin ½ partners. It turns out that one requires an additional doubling of the "light" Higgses, so that one has both $\underline{5}$ and $\underline{\overline{5}}$ supermultiplets H_{α} and $\overline{H}^{\alpha}$.

These are both necessary if one is to give masses to all the quarks and leptons via terms in the superpotential (Dimopoulos & Georgi 1981; Sakai 1982):

$$P \ni m_{2/3}\, \varepsilon^{\alpha\beta\gamma\delta\varepsilon}\, \chi_{\alpha\beta}\, \chi_{\gamma\delta}\, H_\varepsilon + m_{-1/3}\, \psi^\alpha\, \chi_{\alpha\beta}\, \bar{H}^\beta \qquad (26)$$

and we thereby also cancel the chiral anomalies which would otherwise be generated by the fermions in a single Higgs chiral supermultiplet. It is easy enough in principle to arrange the superpotential couplings so that the effective masses of the SU(2) doublet Higgses in the supermultiplets H, $\bar{H}$ vanish:

$$P \ni \lambda_1 \left[\frac{1}{3}\,\mathrm{Tr}(\phi^3) + \frac{m}{2}\,\mathrm{Tr}(\phi^2)\right] + \lambda_2\, \bar{H}(\phi + 3m')H \qquad (27)$$

all one has to do is choose $m = m'$. This condition is unaesthetic and its origin is unclear. However it is technically natural because of the no-renormalization theorem (Wess & Zumino 1974b; Iliopoulos & Zumino 1974; Ferrara et al. 1974). For more details of the construction of susy GUTs, see Ellis et al. (1982d) and papers cited therein.

We now turn to the phenomenology of the minimal susy GUTs. A first remark is that the rate of approach of the gauge couplings is modified (Dimopoulos et al. 1981; Ibáñez & Ross 1981; Einhorn & Jones 1982) thereby increasing the expected value of the grand unification mass scale: in leading loop order

$$\frac{1}{\alpha_3(Q)} - \frac{1}{\alpha_2(Q)} = \left(\frac{11 + (N_H/2)}{12\pi}\right) \ln\left(\frac{Q^2}{m_X^2}\right) \xrightarrow{\text{susy}} \left(\frac{9 + (3N_H/2)}{12\pi}\right) \ln\left(\frac{Q^2}{m_X^2}\right) \qquad (28)$$

The decrease from 11 to 9 is due to gaugino loops partially counteracting gauge boson loops. The increase in sensitivity to light Higgs multiplets is due to the apparition of light shiggs fermions. Since N_H is even in a susy GUT, the Higgs effects are at least six times greater than in minimal conventional GUTs. With $N_H = 2$ one finds when 2-loop effects are included (Einhorn & Jones 1982; Ellis et al. 1982j)

$$m_X = 6(\pm 3?) \times 10^{16} \times \Lambda_{\overline{MS}} \qquad (29)$$

so that $\Lambda_{\overline{MS}} = 150$ MeV corresponds to m_X about 1×10^{16} GeV. The corresponding value of $\sin^2\theta_W$ is

$$\sin^2\theta_W = 0.236 \pm 0.002 \qquad (30)$$

for $\Lambda_{\overline{MS}}$ = 100 to 200 MeV. This is less comfortably close to the best experimental value (2) than was the conventional GUT prediction (1), but judicious adjustment of the susy particle thresholds may dissolve this conflict. If one goes to $N_H = 4$, Equation (28) shows that m_X is decreased, but then $\sin^2\theta_W$ exceeds 0.255 (Einhorn & Jones 1982; Ellis et al. 1982j). Remarkably enough, the value of m_b/m_τ calculated in susy GUTs is numerically essentially identical with that in conventional GUTs:

$$\left.\frac{m_b}{m_\tau}\right|_{\substack{\text{conv}\\ \text{1 loop}}} : \left[\left.\frac{m_b}{m_\tau}\right|_{\substack{\text{susy}\\ \text{1 loop}}} = \left[\frac{\alpha_3(2m_b)}{\alpha_3(m_\tau)}\right]^{12/23}\left[\frac{\alpha_3(m_\tau)}{\alpha_3(m_W)}\right]^{4/7}\left[\frac{\alpha_3(m_W)}{\alpha_{SUM}}\right]^{8/9}\right]$$

$$= 1 : 1.0 \quad . \tag{31}$$

Looking at the increased value (29) of m_X, one might expect the nucleon lifetime to be longer in susy GUTs than in conventional GUTs, if $\tau_B \propto m_X^4$ as before.

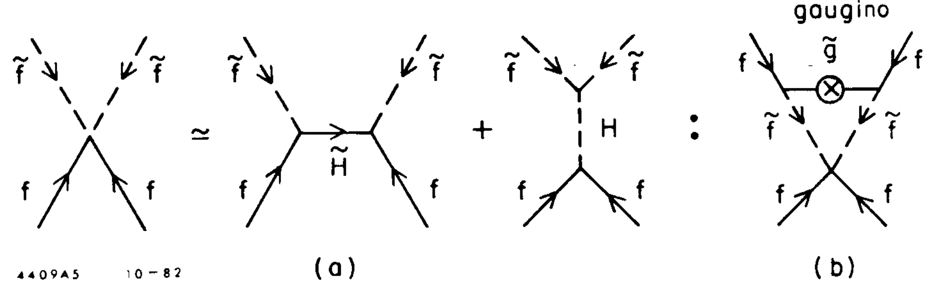
However, Weinberg (1982c) and Sakai & Yanagida (1982) have pointed out that there is a new class of baryon decay diagrams in susy theories, shown in Figure 5, which may give a much shorter lifetime. They involve a dimension 5 $\Delta B = \Delta L$ operator (Figure 5(a)) coupling quarks, squarks, leptons and sleptons arising from Higgs supermultiplet exchange, which has a coefficient $O(1/m_X)$:

$$\frac{g_Y^2}{m_{H_X}}\left(\tilde{q}\,\tilde{q}\,q\,\ell \ \text{ or } \ q\,q\,\tilde{q}\,\ell\right) \quad . \tag{32}$$

The operator (32) does not yet give baryon decay, but must be dressed as in Figure 5(b) by the exchange of a gaugino with a supersymmetry breaking mass which we take to be $m_{\tilde{W}} = O(m_W)$, to give a conventional effective $qqq\ell$ operator. The coefficient of this operator will be

$$O\left(\frac{\alpha}{4\pi}\right)\frac{g_Y^2}{m_{H_X} m_{\tilde{W}}} \tag{33}$$

Fig. 5. (a) Diagrams contributing to the dimension 5 operators responsible for baryon decay in minimal susy GUTs, and (b) the dimension 5 operator dressed by an external gaugino loop (Weinberg 1982c; Sakai & Yanagida 1982).



and we therefore have $\tau_B \propto m_X^2 m_{\tilde{W}}^2$ rather than m_X^4. One might worry that the baryon lifetime may now be too short, but this appears not to be the case (Dimopoulos et al. 1982; Ellis et al. 1982j). If one compares the coefficients of the effective operators for baryon decay in susy and conventional GUTs, one sees that the susy interaction is boosted by a factor $O(m_X/m_W) = O(10^{12})$ but that this is compensated by Yukawa couplings g_Y^2 (32) of $O(10^{-8})$, a loop factor (33) of $O(10^{-2})$, a short distance suppression factor of $O(10^{-1})$, and the change in m_X^{-1} (29) of $O(10^{-1})$. Thus the conventional and susy baryon decay rates are comparable, and a susy baryon has $\tau_B \gtrsim O(10^{30})$ years if

$$m_{\tilde{W}}\, m_X > O(10^{18})\ \mathrm{GeV}^2 \tag{34}$$

or

$$m_{\tilde{W}} \gtrsim O(100)\ \mathrm{GeV} = O(m_W) \tag{35}$$

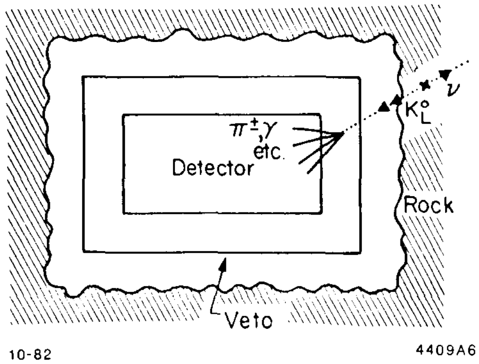
if $m_X = O(10^{16})$ GeV as suggested by Equation (29).

A possible snag, however, is that susy baryon decay modes are very different from conventional baryon decay modes (Dimopoulos et al. 1982; Ellis et al. 1982j):

$$B \rightarrow \bar{\nu}K >> \bar{\nu}\pi >> \mu^+\pi >> e^+K >> e^+\pi \tag{36}$$

and do not match very well the characteristics of the baryon decay candidates reported (Krishnaswamy et al. 1981, 1982) from the Kolar Gold Fields which look more like $B \rightarrow e^+\pi, \rho, \omega$. Many of the proposed new baryon decay experiments are optimized for these decay modes, but would be less efficient at detecting $B \rightarrow \bar{\nu}K$. An idea (Rozanov 1982) for a detector to look for $B \rightarrow \bar{\nu}K$ decays is shown in Figure 6. One looks for $n \rightarrow \bar{\nu}K_L^o$ decays in the rock surrounding a cave, with counters around the walls to veto incoming particles, and a detector inside to pick up any penetrating $K_L^o \rightarrow \pi e\nu$ or $\pi\pi$ or $\pi\pi\pi$ decays, reconstruct them and look for a monochromatic K_L^o signal coming from baryon decay. It seems likely that one could search for susy baryon decays (36) with a sensitivity to lifetime of order 10^{32} years.

Fig. 6. The idea of Rozanov (1982) for detecting $n \to \bar{\nu} K_L^o$ decays.



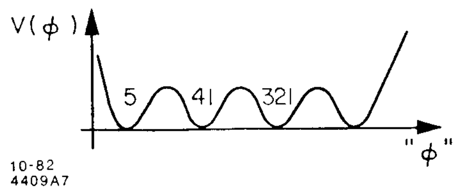
4 SUSY GUTS AND COSMOLOGY

Life in the very early Universe is certainly more complicated and potentially interesting in susy GUTs. We have already seen (22) the form of the effective scalar potential in a susy theory at zero temperature. The potential (22) is clearly positive semi-definite, and is zero if and only if susy is unbroken. Realistic models often have several susy minima which are degenerate with zero vacuum energy, at least in the absence of radiative corrections. For example, the minimal SU(5) model (26,27) discussed earlier has degenerate susy minima which are

$$\left.\begin{array}{ll}
\text{SU(5) invariant} & : \langle 0|\phi|0\rangle = 0 \\
\text{SU(4)} \times \text{U(1) invariant} & : \langle 0|\phi|0\rangle = \frac{V}{2\sqrt{2}}\begin{pmatrix} 1 & & & & 0 \\ & 1 & & & \\ & & 1 & & \\ & & & 1 & \\ 0 & & & & -4 \end{pmatrix} \\
\text{SU(3)} \times \text{SU(2)} \times \text{U(1) invariant} & : \langle 0|\phi|0\rangle = \frac{V}{\sqrt{30}}\begin{pmatrix} 2 & & & & 0 \\ & 2 & & & \\ & & 2 & & \\ & & & -3 & \\ 0 & & & & -3 \end{pmatrix} \\
\text{Dragon (1982) minimum} & : \langle 0|\phi|0\rangle = \frac{V}{\sqrt{30}}\begin{pmatrix} 2 & & & & 0 \\ & 2 & & & \\ & & 2 & A & \\ & & & -3 & \\ 0 & & & & -3 \end{pmatrix}
\end{array}\right\} \quad (37)$$

and more complicated possibilities exist in other susy GUTs (Frampton & Kepkart 1982; Buccella et al. 1982 a,b). The qualitative situation is shown in Figure 7: at high temperatures $T = O(m_X)$ the Universe may not

Fig. 7. Degenerate susy potential minima in minimal susy SU(5) (Dimopoulos & Georgi, 1981; Sakai 1982).



know which of these minima to choose, and different causally separated regions of the very early Universe may choose differently.

There are several possible ways of breaking the degeneracy between the different susy minima. Clearly susy breaking effects do, but they give

$$\delta V \sim m_S^2\, m_X^2 \quad \text{or} \quad m_S^4 \tag{38}$$

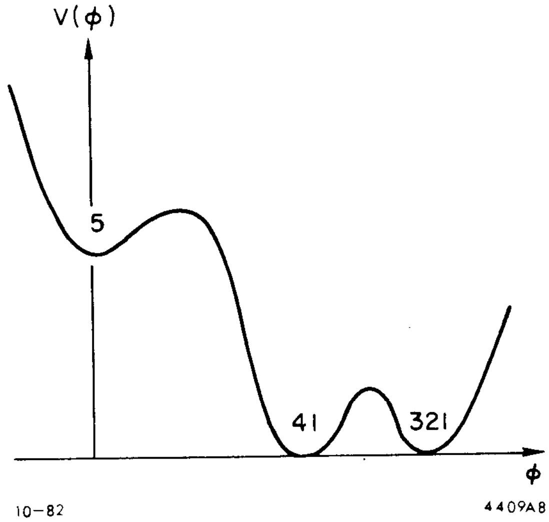
where m_S is a susy breaking parameter which is possibly much less than m_X. Even if the SU(3) × SU(2) × U(1) minimum is favored by these effects (38), the rate at which parts of the Universe which had cooled into one of the other minima (37) then tunnel into the energetically preferred one may be unacceptably slow (Ellis et al. 1982f; Srednicki 1982b):

$$\text{Tunnelling probability} \propto e^{-B} : B = O\left(\frac{m_X}{m_S}\right)^6 \lambda^4 \tag{39}$$

where λ is a generic chiral superfield coupling. Either one wants $m_S \gtrsim m_X$ (which may run into other difficulties) or else one must take $\lambda << 1$ if one is to get an adequately short lifetime for the false minima (Nanopoulos et al. 1982).

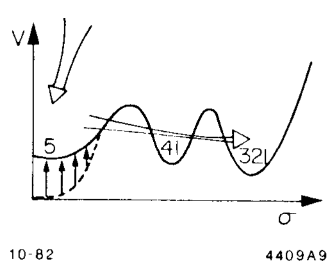
In point of fact, finite temperature fluctuations may get one out of the false minima more quickly than quantum tunnelling. Finite temperature effects give energy shifts proportional to the number of light particles. Unfortunately, in minimal SU(5) this means that the SU(5) invariant minimum (37) is the lowest, and the SU(4) × U(1) and SU(3) × SU(2) × U(1) minima are an equal amount higher. However, it is possible to construct variants of the simplest susy GUT (26,27) in which SU(5) invariant minimum actually breaks supersymmetry, and is hence is higher than the SU(4) × U(1) and SU(3) × SU(2) × U(1) minima (Ellis et al. 1982d), while finite temperature effects favor the desired SU(3) × SU(2) × U(1) minimum as shown in Figure 8.

Fig. 8. Possible form of the potential in non-minimal susy SU(5) (Ellis et al. 1982d).



There is another cute idea for getting out of an undesirable SU(5) minimum should it exist (Nanopoulos & Tamvakis 1982a; Sredincki 1982a). The point is that in the SU(5) invariant phase the gauge coupling tends to become strong at an energy or temperature $O(10^{10})$ GeV. This may lead to the degeneracy between the various minima being broken at this level. For example, when the SU(5) coupling is strong, presumably it confines, and the number of light degrees of freedom may be reduced, so that it becomes energetically disfavored by comparison with the $SU(4) \times U(1)$ and $SU(3) \times SU(2) \times U(1)$ minima. If the generic Yukawa coupling λ is sufficiently small, the Universe may then find its way to the desired $SU(3) \times SU(2) \times U(1)$ minimum as indicated in Figure 9. It may be that the transition passes through an intermediate $SU(4) \times U(1)$ phase with the previous history recycled at a strong coupling temperature $O(10^6)$ GeV.

Fig. 9. Possible non-perturbative SU(5) effect on the minima in a susy GUT (Nanopoulos & Tamvakis 1982a; Srednicki 1982a).



Since SU(5) is broken at temperatures << m_X in this scenario, one must rethink Big Bang baryosynthesis. An attractive option is to have relatively light "heavy" color triplet Higgses with masses $O(10^{10}$ to $10^{11})$ GeV. This is the lowest mass compatible with $d = 6$ Higgs exchange operators letting baryons live longer than 10^{30} years. If dominant, they give (Nanopoulos & Tamvakis 1982b)

$$B \to \mu^+ K' \text{ or } \bar{\nu} K \quad . \tag{40}$$

If there are such light color triplet Higgses, one must find some extra symmetry to forbid them from making a $d = 5$ operator contribution to baryon decay (cf. Figure 5(a)).

Either of the scenarios in Figure 8 or Figure 9 can probably generate sufficient baryon number in the very early Universe.

5 SUPERSYMMETRIC INFLATION

The other currently exciting application of GUTs in the very early Universe is to generate cosmological inflation (Guth 1981; Linde 1982a; Albrecht & Steinhardt 1982). This can be driven by a vacuum potential energy density which is much larger than T^4. *A priori* it might seem to be a disadvantage that in susy theories the vacuum energies at (broken) susy minima are small (38) which might make it more difficult to drive inflation. Furthermore, the smallness of the vacuum energy may make it difficult to reheat to a high enough temperature to achieve baryosynthesis. However, both these difficulties can be resolved if one recalls that susy is automatically broken at finite temperatures, and hence there is no conclusive reason why the Universe should be in a supersymmetric configuration while it is inflating. If susy is broken the vacuum energy can be much larger than suggested by (38). Indeed, susy may even be beneficial in achieving sufficient inflation without generating unacceptably large fluctuations (Ellis *et al*. 1982 g,h,i). This is because susy allows one to adjust or fine-tune parameters without having to worry about them being disturbed by radiative corrections.

To see this, let us first (Ellis *et al*. 1982g) look at the conventional Coleman-Weinberg potential

$$V(\phi) = A\phi^4(\ln \phi^2/\sigma^2 - \tfrac{1}{2}) + D\phi^2 \tag{41}$$

where

$$A = (1/64\pi^2\sigma^4)\left(\sum_B g_B m_B^4 - \sum_F g_F m_F^4\right) \tag{42}$$

with $g_{B(F)}$ the number of boson (fermion) helicity states coupled to the field ϕ. In minimal SU(5) one has

$$A = \frac{5625}{1024\pi^2} g^2 \quad . \tag{43}$$

The parameter D is Equation (41) is an effective $(\text{mass})^2$ parameter which can get contributions from many sources:

$$D = \frac{1}{2}\left(m_0^2 + cT^2 + bR - 3\lambda\langle\phi^2\rangle\right) \tag{44}$$

where m_0 is the flat-space, zero-temperature mass, $c = (75/8)g^2$ and T is probably the Hawking temperature $T_H = H/2\pi$ during an inflationary De Sitter epoch, b is an unknown parameter which is probably O(1) and specifically 1/6 for conformally coupled scalar fields, R is the scalar curvature, $-\lambda/4$ is an effective scalar self-coupling and $\langle\phi^2\rangle$ the quantum expectation value of ϕ^2 (Linde 1982b). To get sufficient inflation we must ensure that D is small:

$$\frac{3H^2}{2D} > 65 \tag{45}$$

where

$$H = \left(\frac{4\pi A\sigma^4}{3m_P^2}\right) \tag{46}$$

is the Hubble parameter during the inflationary epoch. The condition (45) requires a delicate suppression of the D term (44) which seems unnatural in a theory which is not supersymmetric. However, susy allows to "set and forget" m_0 at any value we like, the cT^2 term can be arbitrarily small if the field driving inflation is a gauge singlet, the bR term is absent in a conventionally Weyl-rescaled model of chiral supermultiplets coupled to simple $N=1$ supergravity (Cremmer et al. 1978, 1979, 1982 a,b), and the effective scalar coupling λ can be chosen to respect the Linde (1982b) condition

$$\lambda \leq \frac{\pi^2}{1800} \simeq 5 \times 10^{-3} \tag{47}$$

which is not the case in the conventional GUT model (41) where A (43) is typically O(1/10). Even if we get enough inflation, we must guard against having excessively large fluctuations. In our earlier work (Ellis et al. 1982g) mentioned at this meeting, we took the point of view that one only had a chance of suppressing fluctuations to sufficiently low levels if the Hawking & Moss (1982) action

$$B = \frac{3m_P^4}{8}\left[\frac{1}{V_0} - \frac{1}{V_1}\right] \tag{48}$$

was considerably larger than unity. In this case it might be legitimate to view the flip of an entire De Sitter horizon volume from the origin $\phi = \phi_0$, $V = V_0$ to the local maximum of the potential (41) at $\phi = \phi_1$, $V = V_1$ shown in Figure 10. If B were only of order unity, there would be no strong reason for large regions of space to become hung up at the local maximum, and one would expect the phase transition to be very inhomogeneous and yield unacceptably large fluctuations. The condition that B (48) for the potential (41) be large is:

$$\frac{3m_P^4 D^2}{16A^3\sigma^8} \gg 1 \quad . \tag{49}$$

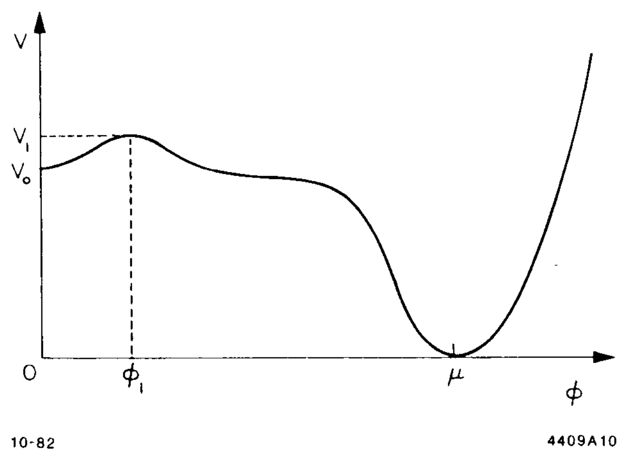
The conditions of sufficient inflation (45) and of smoothness (49) can only be reconciled if A is very small (Ellis et al. 1982g):

$$A < 3(\pi/130)^2 = O(10^{-3}) \tag{50}$$

which is certainly not the case for minimal SU(5) for which A = O(1/10) from Equation (43). Indeed, two recent calculations (Hawking 1982; Guth & Pi 1982) discussed at this meeting suggest that the magnitude of fluctuations in a new inflationary Universe based on minimal SU(5) are indeed much too large.

We have now seen several reasons why susy may help the inflationary Universe: it gives us more freedom to suppress the effective mass parameter D (44) in a natural way, and it enables us to tune down A (42)

Fig. 10. Sketch of the effective potential for ϕ, including the local minimum V_0 at the origin, the barrier provided by the local maximum V_1, and the global minimum $V = 0$ at $\phi = \mu$.



as desired (50) for solving the homogeneity problem. Therefore, since the Workshop we have constructed (Ellis et al. 1982h) a toy supersymmetric model which one can use as a test-bed for supersymmetric inflation. It is defined by a superpotential

$$P(\phi,X,Y) = aX\,\phi(\phi-\mu) + GY(\phi^2-\mu^2) \tag{51}$$

whose associated potential (22,24) has a zero-temperature supersymmetric minimum at $\phi = \mu$, and a non-zero value at $\phi = 0$ which can be used to drive inflation. The conditions (45,49) can be satisfied if

$$a,b \sim O\left(10^{-3}\mu/m_P\right) \quad , \tag{52a}$$

$$\frac{a^2}{b^2} - 2 = O\left(10^{-1}\mu^2/m_P^2\right) \quad . \tag{52b}$$

It is certainly possible to impose these conditions on the supersymmetric model (51), and they are technically natural, but they look like rather inelegant fine-tuning, at least in the case of (52b) when $\mu << m_P$. For this reason, it looks more appealing to postulate that inflation took place close to the Planck epoch, which might entail taking $\mu >> m_X$. In this case we would want the Higgs field driving inflation to be a SU(5) gauge singlet, which is indeed how we chose ϕ in the superpotential (51) and incidentally enables us to suppress the finite temperature contribution to D in Equation (44). The hypothesis of inflation driven by an order parameter much larger than m_X we term "primordial inflation" (Ellis et al. 1982h).

There is another reason for favoring this hypothesis which is related to the fluctuation problem. We have recently used (Ellis et al. 1982i) the formalism of Olson (1976) developed by Guth & Pi (1982) to analyze the spectrum of perturbations expected in the supersymmetric toy model (51). We find that they are approximately scale invariant in a way not very dissimilar from that found by Guth & Pi (1982), and with a magnitude which can be adjusted by varying a and b in the model (51). We get the desired spectrum $\delta\rho/\rho = O(10^{-4})$ if (Ellis et al. 1982i)

$$a,b = O\left(10^{-6}\mu/m_P\right) \quad . \tag{53}$$

This indicates that the previous Hawking & Moss (1982) condition (49) which becomes (52a) for our toy model, while necessary to suppress inhomogeneities, is not in itself sufficient. The fluctuation condition (53) appears to push us in the direction of very small a,b unless $\mu = O(m_P)$. Another motivation for primordial supersymmetric inflation?

One snag of primordial inflation is that it reopens the grand unified monpole problem. Inflation no longer suppresses the monopole density and one must appeal to another mechanism. Two possibilities are the supercosmology (Nanopoulos & Tamvakis 1982 a,b; Sredincki 1982 a,b) scenario outlined in Section 5, and the possibility (Bais & Rudaz 1980) that monopole production may be suppressed by thermodynamic factors $\exp[-(m_M/T_G)]$ during a second order GUT phase transition (Ellis et al. 1982h).

6 IS GRAVITY AN ESSENTIAL COMPLICATION?

So far we have neglected gravity in our discussion of GUTs and susy GUTs, and generally one thinks that this is likely to be a reasonable first approximation. However we have already seen a couple of ways in which gravitational effects may show up in inflationary cosmology. One is the R term in the expression (44) for the effective Higgs $(\text{mass})^2$ D. The other is the T^2 term in (44) which may be expected to become the square of the Hawking temperature $T_H = H/2\pi$ during the De Sitter inflationary epoch. The coefficient b of the R term in (44) is arbitrary until we know how scalar fields couple to gravity, while treating finite curvature effects solely through this and a T_H^2 term may not be completely satisfactory. Another reason why we might worry about gravitational effects is the impetus (52a,53) that we received towards inflation with $\mu = O(m_P)$. Finally, there is the more abstract theoretical point that if one wants to make a theory with local supersymmetry then one must incorporate gravity and construct a supergravity theory, which is indeed the only consistent framework for combining susy and gravity.

As mentioned in Section 2, there are supergravities with $N = 1,2,3,\ldots,8$ extended supersymmetries (Van Nieuwenhuizen 1981). Even if we believe that the ultimate physical theory has $N > 1$, for the reasons of chirality discussed in Section 2 the only possible effective low-energy theory is one based on $N = 1$. In this theory the supermultiplets (18,19) are supplemented by the

$$\text{graviton-gravitino supermultiplet} : \begin{pmatrix} 2 \\ 3/2 \end{pmatrix} . \qquad (54)$$

In addition to the new gravitino particle and its couplings, there are additional low-energy non-renormalizable effective interactions among familiar particles (18,19) which are scaled by universe powers of the

Planck mass (Cremmer et al. 1978, 1979, 1982 a,b). An important example is the effective scalar potential. It is convenient to write the modified potential in a Weyl-rescaled form where there is no $R\phi^2$ term: $b = 0$ in Equation (44). The potential for gauge-singlet fields A (no D term in Equation (22)) is then

$$V = \exp\left(\tfrac{1}{2}\ |A|^2\right)\left[2|F_A + \tfrac{1}{2}A^*P|^2 - 3|P|^2\right] \tag{55}$$

if we make the simplifying assumption of canonical derivative terms $\frac{1}{2}|\partial_\mu A|^2$ in the Lagrangian. (More general forms are possible, but they do not alter qualitatively the subsequent discussion.) Note that we have used "natural" units $\kappa = 8\pi/3m_P^2 \equiv 1$ in writing the expression (55). When we compare and contrast with the globally supersymmetric potential (22) we see the important difference that a negative term is now present. This enables one to cancel the positive vacuum energy otherwise expected in states with susy broken, and even enables us to have states with negative vacuum energy. Weinberg (1982b) has emphasized that the P-dependent terms in (55) can break the degeneracy between different globally supersymmetric minima, but has shown that even if there are minima with negative vacuum energy, a state with zero vacuum energy is stable against decay into them.

Now that we have introduced the gravitino (54), we must ask and answer a few questions about it, notably what is its mass and how many should there be in the Universe? The gravitino acquires a mass by eating the spin $\frac{1}{2}$ goldstino particle associated with the spontaneous breakdown of global supersymmetry, in much the same way that a gauge boson becomes massive by eating a spin zero Goldstone boson (Van Nieuwenhuizen 1981). The gravitino mass is related to the scale of supersymmetry breaking in such a way that

$$m_{\tilde{G}} = \sqrt{m_S} = \exp\left(\tfrac{1}{4}\,|A|^2\right)|P| \tag{56}$$

if the cosmological constant is set to zero by cancelling the positive and negative terms in (55). We see from Equation (55) that there is a contribution

$$\tfrac{1}{2}\ |A|^2|P|^2 \exp\left(\tfrac{1}{2}\ |A|^2\right) \tag{57}$$

to the spin-zero field mass matrix which implies a lower bound (Ellis & Nanopoulos 1982b)

$$m_0 \gtrsim m_{\tilde{G}} \tag{58}$$

on the average of spin-zero particle masses in a chiral supermultiplet.

We know that at least some spin-zero particles must be quite light, notably squarks, sleptons and Higgs bosons. It is always possible (Gaillard et al. 1982) to set up a spontaneously broken global symmetry in such a way that one of the two spin-zero particles in a given chiral supermultiplet is massless, but we do not know of any realistic examples where both of them are massless at the tree level. Furthermore, in many cases the component which is massless at the tree level is only a pseudo-Goldstone boson because the corresponding global symmetry is broken by gauge or other non-gravitational interactions. In this case the (almost) massless spin-zero component can acquire a $(\text{mass})^2$ which is plausibly $O(\alpha)$ times $m_{\tilde{G}}^2$ (Gaillard et al. 1982). Since both spin-zero components of all quark and lepton supermultiplets must have masses $\lesssim O(10^2)$ GeV in order to protect the gauge hierarchy, while the "light" Higgs bosons must also have masses $\lesssim O(10^2)$ GeV in order to break the weak gauge symmetry at the desired scale, the potential existence of a few pseudo-Goldstone bosons does not help. The average masses m_0 of squarks, sleptons and Higgses must all be $\lesssim O(10^2)$ GeV and therefore the constraint (58) on the average spin-zero boson mass tells us that (Ellis & Nanopoulos 1982b)

$$m_{\tilde{G}} \lesssim O(10^2) \text{ GeV} \tag{59}$$

which in turn (56) means that

$$m_S \lesssim O(10^{10}) \text{ GeV} \quad . \tag{60}$$

This range of m_S means that a susy GUT is still essentially supersymmetric on a mass scale $O(10^{16})$ GeV. However, this does not prevent one from achieving inflation, as was exemplified by our toy model (51).

There is a further cosmological complication with the mass of the gravitino which was pointed out by Weinberg (1982a). It is that gravitinos have relatively long lifetimes:

$$\tau_{\tilde{G}} \gtrsim O\left(m_P^2/m_{\tilde{G}}^3\right) \quad . \tag{61}$$

If present in the very early Universe, they may survive embarrassingly long, and either dominate the Universe with their excess mass density, and/or drown it in entropy when they eventually decay (the "gravitino problem"). Weinberg (1982a) suggested that either gravitinos should decay before nucleosynthesis in which case

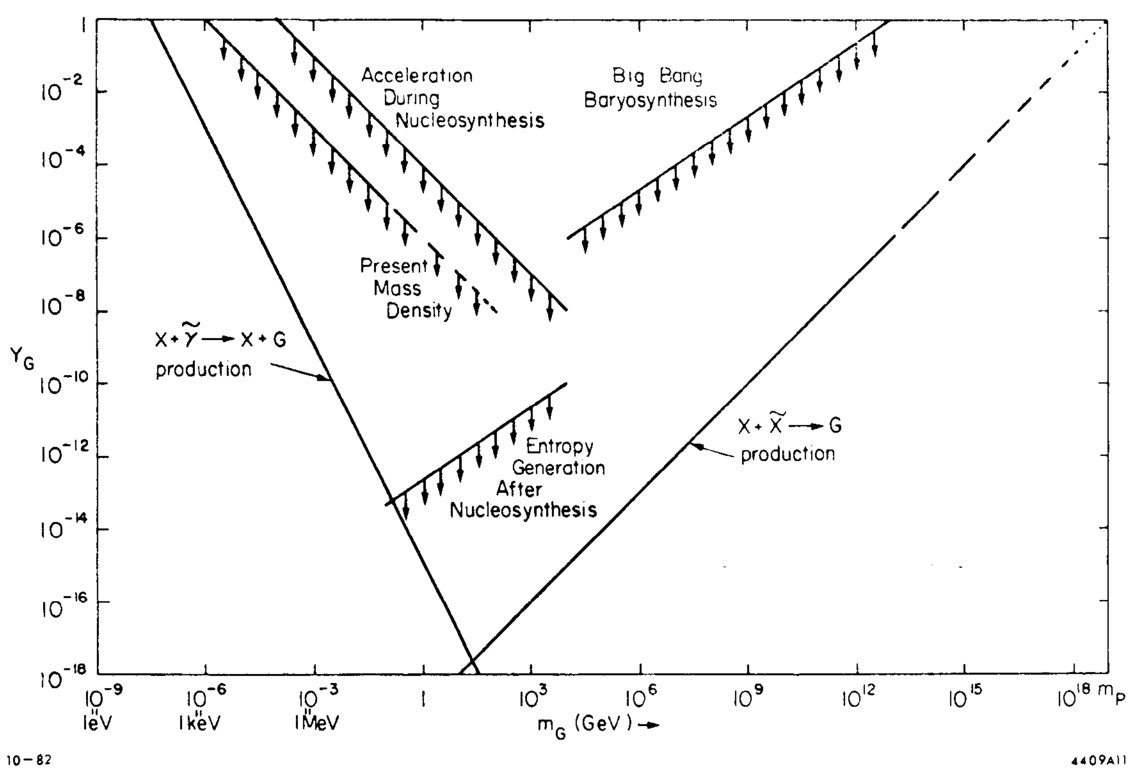
$$m_{\tilde{G}} \gtrsim O(10^4) \text{ GeV} \quad , \quad m_S \gtrsim O(10^{11}) \text{ GeV} \tag{62}$$

or else they should be sufficiently light to avoid contributing too much density to the present-day Universe (Pagels & Primack 1982):

$$m_{\tilde{G}} \lesssim O(1)\text{ keV} \quad , \quad m_S \lesssim O(10^6 \text{ to } 10^7)\text{ GeV} \quad . \tag{63}$$

The lower limit (62) is in prima facie conflict with the upper limit (59): does this mean we have to choose the lower range (63) of gravitino masses? Not necessarily, because we have shown during this Workshop (Ellis et al. 1982a) that even a relatively modest amount of inflation can suppress the primordial gravitino number density sufficiently low to make the conditions (62,63) unnecessary, while gravitino production after the Planck epoch is not sufficient to recreate the gravitino problem. Our results are exhibited graphically in Figure 11, which shows that a suppression factor $Y_G \lesssim O(10^{-15})$ of the gravitino number density relative to that naively expected in an adiabatically expanding Universe would suffice to respect all Big Bang nucleo- and baryosynthesis constraints, while gravitino production after the Planck epoch is not important. The amount of inflation required to solve the gravitino problem is relatively modest, since an increase in the

Fig. 11. Cosmological constraints on the abundance Y_G of gravitinos relative to the normal cosmological abundance. An inflation factor L would give a primordial gravitino abundance $Y_G = O(1/L^3)$. Also shown in the figure are some typical gravitino production mechanisms which respect the cosmological constraints.



cosmic scale factor by $O(10^5)$ suffices to suppress Y_G by $O(10^{-15})$. We therefore think that all gravitino masses in the particle physics range (59) may be reconciled with the cosmological constraints.

It is clear that the next step should be to rework supersymmetric inflation using the supergravity effective potential (55). Also one should investigate more closely the cosmological production of gravitinos around and shortly after the Planck epoch, to see how close to m_P one can push the order parameter of the Higgs field driving inflation (cf. the parameter μ of Equations (51),(52b) and (53)). So far we do not have a complete answer to the question raised in the heading to this section.

7 CONCLUSIONS

One may briefly summarize as follows the status of GUTs and supersymmetric GUTs in the very early Universe as follows. Conventional GUTs do not (yet?) conflict seriously with experiment. Furthermore, they perform baryosynthesis in an elegant way. However, they are embarrassed by the relative paucity of grand unified monopoles which may be explained away by cosmological inflation. Unfortunately, it is difficult to achieve sufficient inflation in a natural way and with sufficiently small fluctuations $\delta\rho/\rho$. Furthermore, conventional GUTs have technical difficulties with maintaining the gauge hierarchy.

These technical problems are resolved in supersymmetric GUTs, whose conventional phenomenology is (almost) equally satisfactory. Unfortunately, susy GUTs are often plagued by a plethora of degenerate vacua which complicate discussion of the GUT phase transition and baryosynthesis. However, susy may alleviate the fine-tuning and fluctuation problems associated with inflation in conventional GUTs. When combined with gravity, susy gives birth to a massive gravitino which presents some cosmological problems that are fortunately soluble.

Life with susy is difficult but potentially rewarding.

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