CERN-TH.4764/87
MAD/TH/87-15
UMN-TH-610/87
MIU-THP-87/16

SUPERCOSMOLOGY REVITALIZED

B.A. Campbell^{*)}, J. Ellis
CERN - Geneva

J.S. Hagelin
Physics Department, Maharishi International University
Fairfield, IA 52556, USA

D.V. Nanopoulos
Physics Department, University of Wisconsin
Madison, WI 53706, USA

and

K.A. Olive
School of Physics and Astronomy, University of Minnesota
Minneapolis - MN 55455, USA

A B S T R A C T

Strong coupling effects can avert excess entropy production in supersymmetric models with a large intermediate scale due to symmetry breaking along a flat direction in the effective potential. This is because strong coupling effects may break supersymmetry giving a large vacuum energy which destabilizes the symmetric field configuration. Alternatively, the reduced number of degrees of freedom in the confined phase destabilize the symmetric field configuration at a high temperature $T \gg m_W$. This principle is exhibited in a simple supersymmetric flipped $SU(5) \times U(1)$ GUT model inspired by the superstring. Compatible scenarios for baryosynthesis are briefly discussed.

^{*)} Permanent address: Physics Department, University of Alberta, Edmonton, Alberta, Canada.

CERN-TH.4764/87
MAD/TH/87-15
UMN-TH-610/87
MIU-THP-87/16
June 1987

Grand unified theories (GUTs)¹⁾ provide attractive solutions to many problems of the standard model, but bring with them some new ones. Prominent among these is the hierarchy problem²⁾. GUTs need at least one high scale m_I of gauge symmetry breaking intermediate between m_W and m_P , and the hierarchies m_I/m_W and $m_P/m_I \gg 1$ must be protected from destruction by radiative corrections. A similar problem arises in models based on the superstring³⁾, in which the gauge symmetry at scales just below m_P tends to be much larger than the $SU(3)_C \times SU(2)_L \times U(1)_Y$ of the standard model. It is generally believed that the hierarchy problem of intermediate scales can only be solved in the context of supersymmetry⁴⁾. Radiative corrections which might destabilize a gauge hierarchy are absent in a supersymmetric theory. Moreover, the effective potential in a supersymmetric theory has many flat directions in field space^{*}, which offer the hope of generating dynamically very different mass scales. Unfortunately, finite temperature effects usually keep the theory in the symmetric phase until $T \sim m_W$, so the dynamical relaxation along one of these flat directions often takes a long time, and is accompanied by excessive entropy release when it finally occurs⁵⁾.

Excess entropy may also be generated by the late decays of metastable sparticles such as gravitinos⁶⁾. Another difficulty is that the additional supersymmetric degrees of freedom offer new ways for baryons to decay, through dimension-5 operators⁷⁾, for instance. Inflation may solve the aforementioned gravitino problem, but only if the reheating temperature $T_R \lesssim 10^{10 \pm 2}$ GeV in the most straightforward scenario where $m_{3/2} \sim m_W$ ⁸⁾. In this case the particles (most probably Higgs triplets H_3) whose decays were responsible for cosmological baryosynthesis would need to have masses $\lesssim 10^{12}$ GeV. These could then lead to excessively rapid baryon decay via dimension-5 operators⁷⁾, whilst dimension-6 operators would not be a problem if $m_{H_3} \gtrsim 10^{11}$ GeV⁹⁾. Nowadays, the conflict between baryosynthesis, entropy generation and baryon stability is less acute, because we have learnt that $m_{3/2}$ may be $\gg m_W$ in non-minimal models, such as the no-scale ones¹⁰⁾ rediscovered in the low-energy theories derived from the superstring. If $m_{3/2} \gg m_W$, there is no excessive entropy release when the gravitino or associated Polonyi, dilaton and axion fields decay¹¹⁾, so the gravitino constraint $T_R \lesssim 10^{12}$ GeV is removed and the Higgs triplets H_3 may be

*) These also carry the danger of many different phases in different parts of the Universe.

heavier, thereby alleviating the problem of baryon stability. But the problem remains of excess entropy generation during dynamical relaxation to an intermediate scale of gauge symmetry breaking along a flat direction in the effective potential⁵⁾.

In this paper, we show how this difficulty is naturally avoided¹²⁾ in a GUT broken down to $SU(3)_C \times SU(2)_L \times U(1)_Y$ by some v.e.v. $\langle \Phi \rangle$ at a high intermediate scale $\langle \Phi \rangle = m_I$. Such a GUT group must be large, and so its coupling in the unbroken phase becomes strong at a high energy scale $\Lambda_C \gg m_W$. Phase transitions in the strong coupling régime T , $\langle \Phi \rangle \lesssim \Lambda_C$ can break supersymmetry giving a large vacuum energy $\sim \Lambda_C^4$ and/or confine many degrees of freedom. Therefore the free energy in this type of strong coupling phase where T , $\langle \Phi \rangle \lesssim \Lambda_C$ is higher than in the weak-coupling phase where $\langle \Phi \rangle > \Lambda_C$. Therefore a hot Universe with $\langle \Phi \rangle = 0$ at $T > \Lambda_C$ will make an adiabatic transition to $\langle \Phi \rangle > \Lambda_C$ when $T < \Lambda_C$ ¹²⁾. The field Φ will then relax to $\langle \Phi \rangle = m_I$ at $T = 0$ without excess entropy generation provided Λ_C is sufficiently large. We illustrate the principles of this strong coupling scenario¹²⁾ with the example of a recently revitalized supersymmetric flipped $SU(5)$ GUT^{13),*)} inspired by the superstring. We also sketch how baryosynthesis may be achieved¹⁴⁾ in this model by the decays of particles with mass $< 10^{15}$ GeV without causing rapid baryon decay via dimension-5 operators.

First we sketch our general scenario and discuss entropy generation. Models with large intermediate-scale symmetry breaking acquire it along a direction Φ in field space which is D-flat and F-flat, with the only field-dependence given by a supersymmetry-breaking term $\sim \tilde{m}^2 \Phi^2$ where $\tilde{m}^2 = 0(m_W^2)$ and is negative for small Φ , and possibly non-renormalizable terms $\sim \Phi^{4+n}/m_P^n$. The v.e.v. $\langle \Phi \rangle \sim m_I$ is of order the scale Φ_0 at which $\tilde{m}^2$ changes sign to become positive for large Φ , or $O(\tilde{m}_P^2 m_P^n)^{1/2+n}$ if there are non-renormalizable terms and this estimate is smaller than Φ_0 . In a weakly-coupled theory, finite-temperature effects $\sim \Phi^2 T^2$ overcome the $\tilde{m}^2 \Phi^2$ term and keep the origin stable until $T \sim |\tilde{m}| \sim m_W$ ⁵⁾. However, in a grand unified theory, the asymptotically-free gauge coupling becomes strong¹²⁾ for temperatures T and values of Φ less than some condensation scale Λ_C . At and below Λ_C the theory is expected to form bilinear condensates, breaking gauge and global symmetries dynamically. We expect that this symmetry breaking will proceed according to a straightforward extension of the Most Attractive Channel

*) This scenario does not work for intermediate-scale models whose non-Abelian gauge group factors are all $\subseteq SU(3)$.

(MAC) hypothesis, which we call stumbling^{16),*)} via SMAC. In some cases, the resulting bilinear condensates break supersymmetry, giving rise to a large vacuum energy $\sim \Lambda_c^4$ in the strong coupling phase of the theory when $T, \Phi \lesssim \Lambda_c$. This strong coupling phase is then energetically disfavoured in comparison with the Higgs phase at large Φ , in which the GUT group is spontaneously broken. Hence strong coupling can lead to an effective contribution to the potential of height $\Delta V \sim \Lambda_c^4$ when $\Phi \lesssim \Lambda_c$ which prevents Φ from being trapped in the symmetric vacuum once $T \lesssim \Lambda_c$.

Alternatively, it is possible that bilinear condensates might form in a way which preserves supersymmetry. In this case the end result of the strong coupling dynamics for $\Phi, T \lesssim \Lambda_c$ is often a phase with fewer light degrees of freedom than in the spontaneously broken Higgs phase where $\Phi > \Lambda_c$ ¹²⁾. In general, there is a negative contribution to the free energy at finite temperature which is proportional to the number of light degrees of freedom with $m \ll T$

$$\Delta F = -\frac{\pi^2}{90} N_{\text{eff}} T^4 : N_{\text{eff}} = N_B + \frac{7}{8} N_F \quad (1)$$

The smaller number of degrees of freedom in the strong coupling phase will cause it to be disfavoured energetically in comparison with the Higgs phase at large Φ ¹²⁾.

Hence both scenarios can prevent the Universe from being trapped in the symmetric phase until $T \lesssim O(\tilde{m}) \sim m_W$ as in theories without strong coupling. In our case, the $\tilde{m}^2 \Phi^2$ term in the effective potential is insignificant for $\Phi \sim \Lambda_c$, and the theory makes a transition to the Higgs phase when $T \sim \Lambda_c$. Very possibly, as we show in some toy examples below, there may be no barrier between the dynamically broken phase and the Higgs phase, in which case Φ just rolls or fluctuates into the new phase. However, we cannot exclude the possibility that the strong dynamics may provide a barrier of height $\Delta V \sim \Lambda_c^4$ when $T \sim \Lambda_c$. Since thermal and strong-coupling fluctuations are both $\sim \Lambda_c$ at this epoch, and since the difference in heights on either side of the putative barrier is $\sim T^4 \sim \Lambda_c^4$, we expect an $O(1)$ probability for completion of the transition to the Higgs phase when $T \sim \Lambda_c$ without any excess entropy generation. Clearly an exact estimate of the transition probability depends on unknown details of the potential in the non-perturbative régime.

*) In the published version of Ref. 16), "stumbling" has inadvertently been rendered as "tumbling", making the paper well-nigh impenetrable.

Given the occurrence of such a strong coupling phase transition, we can then estimate how much entropy will be generated as Φ subsequently evolves from a value $\sim \Lambda_c$ when $T \sim T_c$ to its low-temperature value $\sim m_I$. The avoidance of excessive entropy generation will constrain allowable values of Λ_c and m_I . The evolution of Φ at lower T is governed by the equation

$$\ddot{\Phi} + 3H\dot{\Phi} + \tilde{m}^2\Phi = V'(\Phi) : V'(\Phi) \equiv \frac{\partial V}{\partial \Phi} \quad (2)$$

As in the case of the axion¹⁷⁾, the oscillations which start at $T_i \sim T_c$ with a stored energy density $\rho_\Phi \approx \tilde{m}^2 m_I^2$ evolve adiabatically with an energy density

$$\rho_\Phi \approx \tilde{m}^2 m_I^2 (R_i/R)^3 \quad (3)$$

for $R > R_i$ where R_i is the value of the cosmological scale factor when Φ oscillations begin. These oscillations eventually decay when $H \sim \rho_\Phi^{1/2}/m_P \sim \Gamma \sim \tilde{e} m^3/m_I^2$, at which time the energy density is dominated by ρ_Φ . Decays occur at $R = R_D$ where

$$(R_i/R_D) \approx e^{2/3} \frac{\tilde{m}^{4/3} m_P^{2/3}}{m_I^2} \quad (4)$$

The entropy produced in the Φ decays is then given by

$$S_D \approx [\rho(R_D)]^{3/4} \approx e^{3/2} \tilde{m}^{9/2} m_P^{3/2} / m_I^3 \quad (5)$$

to be compared with the entropy in radiation, which is given by

$$S_i(R_D) = T_i^3 (R_i/R)^3 \quad (5')$$

where T_i is the temperature of the radiation when oscillations begin. Comparing Eqs. (5) and (5'), we see that the entropy produced by the Φ oscillations is

$$\Delta \equiv \frac{S_D}{S_i(R_D)} \approx e^{1/2} m_I^3 \tilde{m}^{1/2} / m_P^{1/2} T_i^3 \quad (6)$$

A conservative upper bound on Δ subsequent to baryosynthesis is $\Delta \lesssim 10^{+6}$, though this depends somewhat on the specific mechanism for generating the baryon asymmetry¹⁴⁾. This corresponds in Eq. (6) to

$$(T_i/m_I) \gtrsim 10^{-2} \left(\frac{\tilde{m}}{\epsilon m_P} \right)^{1/6} \approx 10^{-4.7} e^{-1/6} \quad (7)$$

In the absence of inflation, or if after inflation $T_R > \Lambda_c$, the temperature will then decrease until $T \sim \Lambda_c$, where strong coupling effects destabilize the symmetric field configuration as we have previously discussed. Then, once T drops to a value $T \sim \lambda_c$, where λ is the superpotential coupling which gives a mass $\propto \langle \Phi \rangle$ to the lightest superheavy particle, the Φ field starts oscillating. Hence, one should associate T_i in Eqs. (5) to (7) with $\lambda \Lambda_c$. The same holds even if $\lambda \Lambda_c < T_R < \Lambda_c$ provided either (a) $T_R > \lambda \Phi_R$ (where Φ_R is the Φ field expectation value immediately after inflation), in which case Φ will find itself temporarily at a minimum between Λ_c and T/λ , or (b) $T_R \gtrsim \sqrt{m m_I}$, in which case Φ will fluctuate out of its local minimum at $\Phi = m_I$ and into the global minimum between Λ_c and T/λ , where it will remain temporarily until $T < \lambda \Lambda_c$. Finally, if $T_R < \lambda \Lambda_c$ or if $\lambda \Lambda_c < T_R < \min(\lambda \Phi_R, \sqrt{m m_I})$, the Φ field starts oscillating immediately after inflation and T_i should be identified with T_R . In any case, T_R ($> T_i$) is always subject to the above constraint on T_i . An additional constraint on this type of model is that the temperature of the Universe after Φ decay must be greater than $O(1)$ MeV so that nucleosynthesis occurs during a radiation-dominated epoch, rather than during the Φ oscillations. The temperature of the radiation at $R = R_D$ is just

$$T_r \approx S_D^{1/3} \approx \epsilon^{1/2} m^{3/2} m_P^{1/2} / m_I \quad (8)$$

Thus we obtain the additional constraint

$$m_I \lesssim \epsilon^{1/2} m^{3/2} m_P^{1/2} / O(10) \text{ MeV} \approx \epsilon^{1/2} \cdot 10^{17} \text{ GeV} \quad (9)$$

on any intermediate scale model.

We illustrate the above general analysis with a specific supersymmetric flipped $SU(5) \times U(1)$ GUT¹³⁾ inspired by the superstring. It contains three generations of chiral matter fields with the $SU(5) \times U(1)$ transformation properties

$$F_i \equiv (10, 1), \quad \bar{F}_i \equiv (\bar{5}, -3), \quad l_i^c \equiv (1, 5) \quad (10)$$

and the particle assignments

$$F = \begin{pmatrix} 0 & d^c & -d^c & d & u \\ -d^c & 0 & d^c & d & u \\ d^c & -d^c & 0 & d & u \\ -d & -d & -d & 0 & \nu^c \\ -u & -u & -u & -\nu^c & 0 \end{pmatrix}, \quad \bar{F} = \begin{pmatrix} u^c \\ u^c \\ u^c \\ \nu \\ e \end{pmatrix}, \text{ etc.} \quad (11)$$

for the quarks and leptons of each generation. The minimal version of the model also contains two conjugate pairs of Higgs supermultiplets

$$H \equiv (10, 1), \quad \bar{H} \equiv (\bar{10}, -1); \quad h \equiv (5, -2), \quad \bar{h} \equiv (\bar{5}, 2) \quad (12)$$

as well as four singlets $\phi_m \equiv (1, 0)^*$. The superpotential for the model is

$$W = \lambda_1^{ij} F_i F_j h + \lambda_2^{ij} F_i \bar{F}_j \bar{h} + \lambda_3 \bar{F}_i l_j^c h \\ + \lambda_4 H H h + \lambda_5 \bar{H} \bar{H} \bar{h} \\ + \lambda_6^{im} F_i \bar{H} \phi_m + \lambda_7^m h \bar{h} \phi_m + \lambda_8^{mnp} \phi_m \phi_n \phi_p \quad (13)$$

The GUT symmetry is broken by a v.e.v. along the unique F- and D-flat direction $\langle H \rangle = \langle \bar{H} \rangle \neq 0^{**}$. These v.e.v.'s can be rotated into the v_H^c ($\bar{v}_H^c$) directions, indicating that the GUT breaking is

$$SU(5) \times U(1) \longrightarrow SU(3)_c \times SU(2)_L \times U(1)_Y \\ m_I = \langle H \rangle = \langle \bar{H} \rangle = m_{GUT}$$

The λ_4 and λ_5 superpotential couplings and this v.e.v. pair the d_H^c ($\bar{d}_H^c$) with the D_h ($\bar{D}_h$) in massive Dirac states of mass $\lambda_4(\lambda_5)m_{GUT}$. Likewise, the λ_6 superpotential term and this v.e.v. pair the v_i^c and three of the ϕ_m to form Dirac states of mass $\lambda_6 m_{GUT}$. All the other states have small masses $\sim \langle h, \bar{h} \rangle \sim \tilde{m} \sim 1$ TeV. Thus the Higgs phase of this theory has 45 light quarks and leptons, four light Higgs fields, one singlet from among the ϕ_m , one field Φ which is a linear

*) Inflation or baryosynthesis might motivate the inclusion of more fields.

**) We expect this scale to be of order the supersymmetry-breaking, compactification or string excitation scale, whichever is the smaller. It is not, strictly speaking, an intermediate scale of the type discussed in Ref. 18). The arguments in the third paper of Ref. 18) suggest that it is only possible to generate such a large v.e.v. if supersymmetry-breaking scalar masses are present initially as well as gaugino mass terms.

combination of the v_H^c and $\bar{v}_H^c$, and twelve gauge multiplets, for a total of 63 light supermultiplet degrees of freedom. Note that the decay of the field Φ in our model occurs via the diagrams of Fig. 1, which generate effective D-terms $[v_H^c F^\dagger F]_D$, etc. The decay rate is given by $\Gamma \sim (9\lambda_{1,2,3,7}^4/2048\pi^5) (m_\Phi^2 m_{F,\bar{F},L,\tilde{\Phi}}^2/m_I^2)$, which implies that some of $\lambda_{1,2,3,7}$ and/or $m_{F,\bar{F},L,\tilde{\Phi}}^2$ must be quite large in order to be consistent with the nucleosynthesis constraint (9).

We now analyze the strong-coupling spectrum¹⁹⁾ using the SMAC hypothesis that the condensation of two supermultiplets can be predicted from the quantum number dependence of the single gauge boson exchange potential:

$$V \propto g^2 \Delta C : \Delta C = C_c - C_1 - C_2 \quad (14)$$

where g is the running gauge coupling and C_c , C_1 and C_2 are the quadratic Casimirs of the composite channel and the two elementary supermultiplets, respectively. It has been proposed¹⁵⁾ that a condensate will occur in the composite channel at an energy scale Λ_c where:

$$g^2(\Lambda_c) \Delta C = O(1) \quad (15)$$

This proposal has been supported by lattice gauge theory simulations both in the quenched approximation and with dynamical fermion loops²⁰⁾. In particular, an analysis of SU(2) lattice gauge theory in the quenched approximation with vector-like fermions in the $I = \frac{1}{2}, 1, 3/2$ and 2 representations suggests Casimir scaling of the type (15) with condensation occurring at a scale Λ_c , where:

$$g^2(\Lambda_c) \Delta C \simeq 4 \quad (16)$$

Whilst the critical value of $g^2 \Delta C$ at which dynamical symmetry breaking occurs could depend on the gauge group, and there could easily be violations of Casimir scaling (15) for greatly differing values of C , we will take (16) as our criterion for stumbling.

According to our SMAC hypothesis, the first condensate to form would be the SU(5) gaugino condensate $24 \times 24 \rightarrow 1$, for which $\Delta C = 10$. Such a gaugino condensate would signal supersymmetry breaking, leading to a positive vacuum energy $\sim \Lambda_c^4$. Since this vacuum energy does not fall with temperature as T^4 like V_{eff} , any such

supersymmetry breaking would rapidly destabilize the symmetric field configuration causing a transition to the Higgs phase, as we have previously emphasized.

However, it is possible that this naïve application of SMAC is incorrect, and that supersymmetry may forbid such gaugino condensation, as is known to occur in certain theories²¹⁾. Therefore we now present an alternative stumbling scenario assuming unbroken supersymmetry^{*)}. When we supplement our SMAC hypothesis with this additional assumption, we find that the first condensate to form is $\underline{10}_1 \times \overline{10}_H \rightarrow \underline{1}_1$ ($i=1,2,3,H$), which has $\Delta C = 36/5$. This condensate breaks the full flavour chiral symmetry^{**)}

$$U(4)_{F,H} \times U(1)_{\overline{H}} \times U(3)_{\overline{F}} \times U(1)_{\overline{H}} \times U(1)_{\overline{H}} \times U(1)_R / U(1) \times U(1)_A \quad (17)$$

to $U(3)_1 \times U(1) \times \dots$ and also breaks the R symmetry, giving rise to eight Goldstone bosons which sit in the four bound state chiral supermultiplets $\underline{1}_1$. The next SMAC is $\underline{5}_H \times \overline{5}_{\overline{H}} \rightarrow \underline{1}$, which has $\Delta C = 24/5$ and breaks $U(1)_{\overline{H}} \times U(1)_{\overline{H}}$ in (17) down to $U(1)$. This results in one massless Goldstone boson, which sits in the bound state chiral supermultiplet $\underline{1}$ together with its real supersymmetric partner, which is not a Goldstone boson of any broken compact symmetry and therefore corresponds to a non-compact degeneracy of the vacuum²²⁾. All of the above Goldstone bosons from the $\underline{1}_1$ and the $\underline{1}$ acquire masses when Φ acquires a v.e.v., since they each contain a massive D or d_H^C constituent. This will be important when we consider the detailed shape of $V_{\text{eff}}(\Phi, T)$ in the strong coupling phase of the theory.

Since the $\underline{1}_1$ and $\underline{1}$ condensates give masses to all components of the $\underline{10}$, $\overline{10}$, $\underline{5}$ and $\overline{5}$ representations, the chiral symmetry associated with the remaining light gauge non-singlet fields is just $U(3)_F \times U(3)_{\overline{F}} / U(1) \times U(1)_A$. According to the naïve

*) We still anticipate the formation of colour-singlet bound states containing gauge bosons and/or gauginos. Although the natural mass-scale for such bound states is Λ_c , the presence of an R symmetry would imply that chiral supermultiplets made out of gauginos were massless. However, since R symmetry is broken by $\underline{10}_1 \times \overline{10}_H \rightarrow \underline{1}_1$ condensates, we do not expect these states to be massless in the confined spectrum.

**) This chiral symmetry excludes SU(5)-singlet fields, which do not form condensates. The U(1) factors in the numerator of (17) include one anomalous U(1) factor and one gauged U(1) factor, which are therefore divided out.

SMAC hypothesis, the next condensate to form should be $\underline{10}_i \times \underline{10}_j \rightarrow \bar{\underline{5}}_{ij}$, which also has $\Delta C = 24/5$. Such a condensate would necessarily break supersymmetry by a D-term. If this were present, the vacuum energy would be $\sim \Lambda_c^4$ and the symmetric vacuum immediately destabilized as discussed earlier. However, if we continue to assume that supersymmetry overrules SMAC, we must assume that such $\bar{\underline{5}}_{ij}$ condensates do not form until the next SMAC condensate $\underline{10}_i \times \bar{\underline{5}}_j = \underline{5}_{ij}$ with $\Delta C = 18/5$ also forms, yielding a combined D-flat direction. If they are all aligned, as is energetically preferred, these $\underline{5}_{ij}$, $\bar{\underline{5}}_{ij}$ condensates break the gauge symmetry to $SU(4) \times U(1)$, and also break all the remaining chiral symmetries. Since the required Goldstone bosons are among the $\underline{5}_{ij}$ and $\bar{\underline{5}}_{ij}$ bound state superfields, it is reasonable and conservative to assume that all the $\underline{5}_{ij}$ and $\bar{\underline{5}}_{ij}$ bound states form. Taking account of the Higgs mechanism these bound states include five massless $\underline{4}$ representations of $SU(4)$ from among the $\bar{\underline{5}}_{ij}$ [which transform as $(\underline{6}, \underline{1})$ under $SU(3)_F \times SU(3)_{\bar{F}}$], eight $\underline{4}$'s from among the $\underline{5}_{ij}$ [which transform as $(\underline{3}, \underline{3})$ under $SU(3)_F \times SU(3)_{\bar{F}}$], and 14 $SU(4)$ -singlets $\underline{1}_4$, 9 of which contain a v^c and hence receive masses $\propto \langle \Phi \rangle$. All the components in the $\underline{10}_i$ and $\bar{\underline{5}}_j$ participate in these condensates and thus get masses, except for the 3 $SU(4)$ singlets $\underline{1}_j$ within the $\bar{\underline{5}}_j$, which must be counted among the massless states in the strong coupling phase.

Finally, the bound state $\underline{4}$ and $\underline{4}$ representations are expected to form immediately $\underline{4}_m \times \underline{4}_n \rightarrow \underline{1}_{mn}$ condensates, since $\Delta C = 15/4$ which is larger than $18/5$. All of the possible 40 $\underline{1}_{mn}$ bound states are expected to have masses $\sim \Lambda_c$, since there are no gauge or chiral symmetries to forbid such mass terms from appearing in the superpotential.

The total number of light states in the strong coupling phase is therefore expected to be at most $4(\underline{1}_i) + 1(\underline{1}) + 14(\underline{1}_4) + 3(\underline{1}_j) + 4\phi_m + 3(\chi_1^c) = 29$, of which 17 get masses $\propto \Phi$ and 12 do not. As expected, this is fewer than in the Higgs phase, where there are 107 total superfield degrees of freedom, of which 44 acquire masses $\propto \langle \Phi \rangle$ and 63 do not.

To estimate when the bulk of this dynamical symmetry breaking occurs, we use the middle-of-the-road estimate $\Delta C = 24/5$ which, when substituted into (16) corresponds to $\alpha \equiv g^2/4\pi = 0.065$. A renormalization group analysis of the low-energy effective theory yields the estimate $\alpha_{GUT} = 0.043$. Using the leading order formula

$$\alpha(\mu) \simeq \frac{\alpha_{\text{GUT}}}{1 - \frac{5}{2\pi} \alpha_{\text{GUT}} \ln(m_{\text{GUT}}/\mu)} \quad (18)$$

we calculate

$$\Lambda_c \simeq 2 \times 10^{-4} m_{\text{GUT}} \quad (19)$$

which is consistent with the lower limit (7) based on the entropy generation constraint. Clearly our understanding of strong-coupling dynamics is incomplete, and the estimate (19) of Λ_c could be off by orders of magnitude in either direction^{*}, but at least hope is still alive.

We have made a computer simulation of the evolution of the effective potential when $T \sim \Lambda_c$ which is shown in Fig. 2. To model the Φ dependence of the effective potential when $\Phi, T \sim \Lambda_c$ we use a simple monotonic interpolating function:

$$\exp \left\{ - [(\Phi^2 + T^2)/\Lambda_c^2]^p \right\} \quad (20)$$

to match the finite-temperature effective potential:

$$V_{\text{eff}}(\Phi, T) = \frac{T^4}{2\pi^2} \sum_i (-1)^F \int_0^\infty dx x^2 \ln \left\{ 1 - (-1)^F \exp \left[- (x^2 + m_i^2(\Phi)/T^2)^{1/2} \right] \right\} \quad (21)$$

appropriate to the strong coupling ($\Phi, T \lesssim \Lambda_c$) and Higgs ($\Phi > \Lambda_c$) phases of the theory. The coefficient p in Eq. (21) is an unknown and model-dependent power which is presumably $O(1)$. In evaluating (21) we have taken $\lambda_{4,5,6} = O(1)$, as seems appropriate in the strong-coupling domain. Figure 2a shows the effective potential for $T, \Phi \sim \Lambda_c$ if we take $p = 1$, whilst Fig. 2b is for $p = 2$. We see that there is no barrier to trap Φ close to the origin. Although it might be possible by varying parameters to form an apparent barrier of height $O(\Lambda_c^4)$ when $T = O(\Lambda_c)$, we expect, as was explained above, that even in this case strong-coupling and thermal effects would easily make an end run around any such barrier when $T \sim \Lambda_c$.

^{*}) Indeed, it has been argued²³⁾ that supersymmetric GUTs may effectively be in a strong-coupling régime at m_{GUT} .

To make this scenario complete we must still add cosmological inflation and baryosynthesis. We expect that inflation is produced by additional $SU(5) \times U(1)$ singlet fields ϕ . If we dissociate this model from superstring theories, then we are free to use existing acceptable models of inflation²⁴⁾ based on $SU(N,1)$ supergravity. These models, however, contain either linear or quadratic superpotential couplings and as such cannot be used in the context of the superstring, where all couplings must be at least trilinear.

The production of a net baryon asymmetry is much less problematic in this flipped $SU(5) \times U(1)$ model¹³⁾. The magnitude of the asymmetry will, however, depend on the amount of reheating during inflation. A detailed treatment of baryosynthesis will be given elsewhere¹⁴⁾ and here we simply outline the basic mechanism. In the absence of strong reheating: $T_R < \lambda \Lambda_c$, where λ is the Yukawa coupling which gives mass m_ϕ to the lightest particle capable of generating a baryon asymmetry, we must rely on the decay of the inflaton ϕ to ϕ whose subsequent decay will lead to a baryon asymmetry by its out-of-equilibrium decay. The resulting asymmetry is approximately given by

$$\frac{n_B}{S} \simeq \epsilon' T_R / m_{H \text{ or } \phi} \quad (22)$$

where ϵ' is the net baryon number per decay. On the other hand, if $T_R > \Lambda_c$, thermal equilibrium is restored and we must depend on a small decay rate of either H or ϕ_m . For example, if we consider the decay of H to the combination $FFF\bar{f}$, we can estimate the baryon asymmetry to be

$$\frac{n_B}{S} \simeq \epsilon' \left(\frac{m_H^{5/2} m_P^{1/2}}{m_h^3} \right) \quad (23)$$

and clearly an asymmetry of order 10^{-11} is possible.

We have seen in this paper that the entropy problems of intermediate-scale models can be avoided by strong-coupling effects, in an action replay of the previous "supercosmology" scenario¹²⁾ in conventional supersymmetric $SU(5)$ GUTs. This scenario cannot rescue intermediate-scale models with small low-energy group factors $\subseteq SU(3)$, but is applicable to a recently proposed¹³⁾ flipped supersymmetric $SU(5) \times U(1)$ GUT.

ACKNOWLEDGEMENTS

We would like to thank L. Corners, A. Deans, S. Kelley, R. Parker, R. Ticciati and G. Veneziano for valuable conversations. B.A.C. thanks the CERN Theory Division for its kind hospitality. His research was supported by NSERC Canada under grant A8863. The work of D.V.N. was supported in part by DOE grant DE-AC02-76ER0081 and in part by the University of Wisconsin Research Committee with funds granted by the Wisconsin Alumni Research Foundation. The work of K.A.O. was supported in part by DOE grant DE-AC02-83ER-40105.

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FIGURE CAPTIONS

- Fig. 1 Diagram contributing to a D-term generating a decay rate $\propto \tilde{m}^3/m_I^2$ for the Φ field in the flipped $SU(5) \times U(1)$ GUT model.
- Fig. 2 Sample movies of the evolution with temperature of the effective potential $V_{\text{eff}}(\Phi, T)$ in strongly-coupled $SU(5) \times U(1)$ when $T/\Lambda_c = 0(1)$.

- Figure 1 -

- Figure 2a -

- Figure 2b -