CERN-TH.4613/86
MAD/TH/86-29
UMN-TH-591/86

COMMENTS ON INTERMEDIATE-SCALE MODELS

J. Ellis
CERN, Geneva

K. Enqvist and D.V. Nanopoulos
Department of Physics, University of Wisconsin
Madison, Wis.

K. Olive
School of Physics and Astronomy,
Univ. of Minnesota, Minneapolis, Minn.

ABSTRACT

Some superstring-inspired models employ intermediate scales m_I of gauge symmetry breaking. Such scales should exceed 10^{16} GeV in order to avoid *prima facie* problems with baryon decay through heavy particles and non-perturbative behaviour of the gauge couplings above m_I . However, the intermediate-scale phase transition does not occur until the temperature of the Universe falls below $O(m_W)$, after which an enormous excess of entropy is generated. Moreover, gauge symmetry breaking by renormalization group-improved radiative corrections is inapplicable because the symmetry-breaking field has no renormalizable interactions at scales below m_I . We also comment on the danger of baryon and lepton number violation in the effective low-energy theory.

CERN-TH.4613/86
December 1986

Given that the Theory of Everything (TOE) is some string theory formulated in more than four dimensions at energies $O(10^{19})$ GeV, how do we make contact with physics in four dimensions at energies $\lesssim 1$ TeV? Many compactification schemes lead to four-dimensional gauge groups which are larger than the $SU(3)_c \times SU(2)_L \times U(1)_Y$ of the Standard Model [1], and require another stage of gauge symmetry breaking [2]. This might be achieved by giving intermediate-scale vacuum expectation values of order m_I to one or more $SU(3)_c \times SU(2)_L \times U(1)_Y$ -invariant scalar fields. Candidates in the **27** representations of E_6 which are left over as light fields after Calabi–Yau compactification of the heterotic superstring include the conjugate sneutrino $\tilde{\nu}^c$ and the $SO(10)$ -singlet field N . We present here some general comments on the cosmology and particle physics phenomenology of intermediate-scale models, illustrated, where appropriate, by examples where $\langle 0|\tilde{\nu}^c$ and/or $|N\rangle = O(m_I)$.

Our first comment is that in many such models baryon decay can be mediated by dimension-five operators due to the exchange of particles weighing $O(m_I)$, and that experiments [3] on baryon stability strongly suggest $m_I \gtrsim 10^{16}$ GeV [4]. Our second comment is that the evolution of the gauge couplings α_i at energies $> m_I$ is so rapid in models with several mirror generations that if $\alpha_i \simeq 1$ at $m_P/\sqrt{8\pi} = 2.4 \times 10^{18}$ GeV, m_I should generally be within two orders of magnitude of this scale. Our third comment is that, even in a cold Universe Affleck–Dine [5] scenario [6] for baryosynthesis, m_I should be $\lesssim 10^{10}$ GeV if excess entropy is not to be generated when the intermediate-scale phase transition finally occurs at some temperature $T \simeq m_w$. Our fourth comment is that the usual scenario of weak gauge symmetry breaking via radiative corrections, improved using the renormalization group [7], is inapplicable to the intermediate-scale symmetry breaking, since the field with expectation value $O(m_I)$ has no renormalizable interactions at lower energy scales. We also point out that in generic models with both $\langle 0|\tilde{\nu}^c|0\rangle$ and $\langle 0|N|0\rangle \neq 0$ [8], the effective low-energy theory will have both ΔB and $\Delta L \neq 0$ transitions, the latter in particular because the same (unknown) mechanism for generating an $H_u H_d$ mixing term between the conventional supersymmetric Higgses will also generate Higgs–lepton $H_u L$ mixing. Details are given in a separate publication [9].

We start with a discussion of baryon decay via dimension-five operators due to heavy particles weighing M [10]. These give rise to four-fermion operators with coefficients [4]

$$A = O\left(\frac{\lambda^2}{16\pi^2 M m_w}\right) f(\tilde{m}^2/m_w^2) \quad (1)$$

where the λ are Yukawa couplings, and f depends on the ratios of light sparticle masses $\tilde{m}$ to each other and to m_w . Using the small Yukawa couplings λ characteristic of the first two generations of known fermions, some of us have found [4] that for generic values of $\tilde{m}$ the present limits on baryon decay required

$$M \gtrsim 10^{16} \text{ GeV} \quad (2)$$

with 10^{17} GeV being preferred for complete safety. In models with intermediate scales m_I there are many heavy particles weighing $M = O(\lambda m_I)$, where λ is some possibly small Yukawa coupling, and others weighing $M = O(m_I^2/m_P)$ which could mediate baryon decay. We therefore consider

$$m_I \gtrsim 10^{16} \text{ GeV} \quad (3)$$

as a conservative lower limit on m_I in such models, and would feel safer with $m_I \gtrsim 10^{17}$ GeV. We do not consider $m_I = 10^{14}$ GeV to be acceptable.

Next we consider the renormalization group evolution of the strong gauge coupling above m_I , given by

$$\alpha_s(Q) \simeq \frac{12\pi}{(27 - \frac{3}{2}N_3) \ln(Q^2/\Lambda^2)} \quad (4)$$

where N_3 is the number of colour-triplet chiral superfields, and Λ is the scale at which α_s becomes non-perturbative. Including just the supersymmetric Standard Model particles, $N_3 = 12$, and neglecting the possible contributions of particles weighing λm_I : $\lambda \ll 1$, or m_I^2/m_P , to the evolution of α_s below $Q = m_I$, we expect $\alpha_s(10^{14} \text{ GeV}, 10^{16} \text{ GeV}) = 1/23, 1/24$, and above m_I

$$\frac{\Lambda}{m_I} \simeq \exp \frac{23.6\pi}{(\frac{3}{2}N_3 - 27)} \quad (5)$$

where N_3 now includes all the massive states. In the model of Ref. [8], $N_3 = 66$ and $\Lambda/m_I \approx 400$. If Λ is to be equated with $m_P/\sqrt{8\pi} = 2.4 \times 10^{18}$ GeV, we need $m_I \approx 10^{16}$ GeV. Alternatively, if we were to keep $\Lambda \approx m_P/\sqrt{8\pi}$ and insist on $m_I \approx 10^{14}$ GeV, we would find that $\alpha_s(m_I) \approx 1/38$, which is too small even to accommodate the contributions of the minimal supersymmetric Standard Model particles to the evolution of α_s at scales $m_W < Q < m_I$. Clearly the rate of increase of α_s above m_I is model-dependent, but this is a problem to watch.

The provisional conclusion of the above two paragraphs is that any intermediate scale m_I should exceed 10^{16} GeV.

Next we discuss a cosmological problem for intermediate-scale models: in general, the phase transition which generates an intermediate scale via a large vacuum expectation value along some flat direction in the potential will generate a large amount of entropy which washes out any pre-existing baryon asymmetry [11, 12]. This difficulty has previously been discussed in the context of models with baryogenesis at high temperatures [11, 12]. After repeating that argument here, we show that it can be extended to include models [6] where baryogenesis occurs at $T \approx m_W$. If we denote by N the generic ‘flat’ scalar field with potential

$$V(N, T) = (-|\tilde{m}_N^2| + cT^2)N^2 + \frac{N^{4+n}}{m_P^n} \quad (6)$$

where $|\tilde{m}_N^2| \approx \tilde{m}^2 \approx m_W^2$ is the scale of supersymmetry breaking in the observable sector, then the critical temperature for this transition is $T_c \approx \tilde{m}$ (tunnelling through the barrier at finite

temperature is negligible). Because $m_P \gg \tilde{m}$, the potential (6) appears very flat and N evolves at lower T according to the familiar equation of motion:

$$\ddot{N} + 3H\dot{N} + \tilde{m}^2 N = V'(N) \quad (7)$$

where H is the expansion rate of the Universe. The N field oscillates about the minimum at $\langle 0|N|0 \rangle \equiv m_I \approx (\tilde{m}^2 m_P^4)^{1/(2+n)}$, and the density stored in the oscillations is

$$\rho_N \approx \tilde{m}^2 \hat{N}^2 \quad (8)$$

where $\hat{N} \equiv N - \langle 0|N|0 \rangle$. The oscillations finally decay when $H \approx \Gamma_N$, where the decay rate $\Gamma_N = K^2 \tilde{m}^3/m_I^2$ with $K \approx 10^{-2}$, i.e. when

$$\hat{N} = N_D \sim K^2 \tilde{m}^2 m_P / m_I^2 \quad (9)$$

where we have assumed that the energy density of the Universe is dominated by $\hat{N}$ oscillations so that $H \approx \tilde{m}\hat{N}/m_P$. If we define R_N to be the value of the Universe's scale factor when the amplitude of oscillations $\hat{N} \approx m_I$, then the energy density during the oscillations is

$$\rho_N \approx \tilde{m}^2 m_I^2 (R_N/R)^3 \quad (10)$$

whereas the entropy in radiation is

$$S_i \approx \tilde{m}^3 (R_N/R)^3 \quad (11)$$

Since the oscillations begin when $T \approx \tilde{m}$, the scale factor at $\hat{N} = N_D$ is

$$\frac{R_{DN}}{R_N} \approx m_I^2 / K^{4/3} m_P^{2/3} \tilde{m}^{4/3} \quad (12)$$

so that the entropy after decay is

$$S_D = [\rho_N(R_{DN})]^{3/4} \approx K^3 \tilde{m}^{9/2} m_P^{3/2} / m_I^3 \quad (13)$$

Comparing Eq. (13) with Eq. (11) we find

$$\Delta \equiv \frac{S_D}{S_i(R_D)} \approx m_I^3 / K \tilde{m}^{5/2} m_P^{1/2} \quad (14)$$

and if we want to keep $\Delta < 10^6$ in order to preserve the baryon asymmetry, we must choose $m_I \lesssim 10^7$ GeV [11, 12].

One might have hoped to avoid this limit by postulating [13] that the intermediate-scale phase transition occurs during inflation when the temperature T falls below $\tilde{m}$. Then, if the energy density ρ_R fed into particles after inflation were low enough [$\rho_R < (\tilde{m}^2 m_I^2)$] so as not to restore the original gauge symmetry, the baryon asymmetry would not be washed out. It should, however, be borne in mind that the Hawking temperature provides, in some sense, a lower bound on the effective temperature during inflation. Indeed, there is a likely source of even larger, quantum-mechanical, non-thermal fluctuations in the fields. Because of the flatness of the potential $V(N)$, de Sitter fluctuations [14] would prevent the intermediate-scale field N from settling at the minimum of V . Instead, N will be very far from its minimum and the entropy produced in the subsequent rollover could still wash out the baryon asymmetry. For a reheating temperature $T_R < (\tilde{m} m_I)^{1/2}$, oscillations effectively begin when the scale factor is R_R , which is the value of R at $T = T_R$. Thereafter, the entropy in radiation is

$$S_i \approx T_R^3 (R_R/R)^3 \quad (15)$$

The energy density in the N -oscillations can be taken from eq. (10) with the substitution $R_N \rightarrow R_R$. This assumes, however, that $\tilde{N}_0$ (the initial value of $\tilde{N}$) is $O(m_I)$. Depending on the details of $V(N)$, N_0 may in fact be larger than m_I , thus using eq. (10) will yield a conservative lower bound on the entropy generated. In this case, apart from the replacement $R_N \rightarrow R_R$, eqs. (12) and (13) remain unchanged, while the entropy increase becomes

$$\Delta \approx \tilde{m}^{1/2} m_I^3 / k m_p^{1/2} T_R^3 \quad (16)$$

Since we have assumed $T_R \leq (\tilde{m} m_I)^{1/2}$ the smallest value of Δ occurs when $T_R = (\tilde{m} m_I)^{1/2}$ so that $\Delta \geq \Delta_{\min} = m_I^{3/2} / k m_p^{1/2} \tilde{m}$ and $\Delta_{\min} < 10^6$ only if $m_I < 10^{10}$ GeV. Although this limit is softer than the one of $m_I < 10^7$ GeV previously quoted, it still excludes $m_I \geq 10^{14}$ GeV.

One might have hoped that this kind of limit could be avoided if the baryon asymmetry were produced at low temperatures $T \lesssim \tilde{m}$ after this intermediate-scale phase transition. An example of such a scenario [6] could be the Affleck–Dine [5] mechanism for producing the baryon asymmetry from the decays of coherent oscillations of squarks and sleptons ϕ . However, we will now show that, if one follows the evolution of these fields ϕ in conjunction with N in this scenario, the above limit on m_I is only partially relaxed as seen in the figure.

As in ref. 6, we will consider the effects of the two oscillating fields ϕ and N_0 . Depending on the values of T_R and m_I , we can distinguish several possibilities. First of all, if $T_R > (\tilde{m} \phi_0)^{1/2}$, the energy density is dominated by radiation ($T_R^4 > \rho_\phi = \tilde{m}^2 \phi_0^2$) where ϕ_0 is the initial value of ϕ after inflation. In this case, it is no longer appropriate to discuss the Affleck–Dine mechanism as a scenario for baryosynthesis. If we have $(\tilde{m} m_I)^{1/2} < T_R < (\tilde{m} \phi_0)^{1/2}$ then, as we discussed previously, the N -symmetry is restored, we can write ρ_N as in eq. (10) and

$$\rho_\phi \simeq \tilde{m}\phi^2 \simeq \tilde{m}^2\phi_0^2 (R_R/R)^3 \quad (17)$$

where again R_R is the value of R when $T = T_R$ and R_N is the value of R when N oscillations begin (at $T \sim \tilde{m}$), so that $R_R/R_N = \tilde{m}/T_R$. When the N oscillations begin

$$\rho_N/\rho_\phi \simeq m_I^2 T_R^3 / \phi_0^2 \tilde{m}^3 \quad (18)$$

and taking the minimal value for $T_R = (\tilde{m}m_I)^{1/2}$ we see that $\rho_N/\rho_\phi > 1$ when $m_I > (\tilde{m}^3 m_P^4)^{1/7} \sim O(10^{12})$ GeV. The decays of the N and ϕ fields can be determined by setting the decay rates Γ_N and $\Gamma_\phi \sim \tilde{m}^3/\phi^2$ equal to the expansion rate $H \sim \tilde{m}m_I (R_N/R)^{3/2}/m_P$ so that $R_{D\phi}$, the value of R at ϕ decay, is

$$\frac{R_{D\phi}}{R_N} \simeq \tilde{m}^{2/9} m_I^{2/9} \phi_0^{4/9} / T_R^{2/3} m_P^{2/9} \quad (19)$$

and R_{DN} is given by eq. (12). In this case $R_{DN} > R_{D\phi}$ so we compute the baryon asymmetry at R_{DN} ,

$$n_B \simeq \tilde{m}\phi_0^2 (R_R/R_{DN})^3 \epsilon \phi_0^2 / m_G^2 \quad (20)$$

where ϵ represents the CP violation in ϕ decays and m_G is the GUT scale. Using eq. (12), we find $n_B = k^4 \epsilon \phi_0^4 \tilde{m}^8 m_P^2 / m_I^6 T_R^3 m_G^2$. At R_{DN} , s_D is given by eq. (13) so that

$$\frac{n_B}{s_D} \simeq k \epsilon \phi_0^4 \tilde{m}^{7/2} m_P^{1/2} / m_I^3 T_R^3 m_G^2 \quad (21)$$

Taking representative values $k \sim 10^{-2}$, $\epsilon \sim 10^{-3}$, $(m_G/m_P) \sim 10^{-1}$ and $\phi_0^2 \sim 10^{-7} m_P^2$, the constraint $n_B/s > 10^{-11}$ requires $m_I T_R \lesssim O(10^{16} \text{ GeV}^2)$ in contradiction with $m_I > O(10^{12})$ GeV (see the figure). Actually, one can show that even for lower values of m_I the constraint $m_I T_R < O(10^{16} \text{ GeV}^2)$ holds as long as $m_I \gtrsim O(10^{10})$ GeV. This is because even though ρ_N is not dominant at $R_{D\phi}$, because the energy density due to ϕ decay falls off as R^{-4} , ρ_N becomes dominant before R_{DN} . For $m_I < O(10^{10})$ GeV, ρ_N is never dominant and the standard result for n_B/s of Affleck and Dine [5] is obtained, as seen in the figure.

The final case which is of greatest interest here is the case where $T_R < (\tilde{m}m_I)^{1/2}$. The N -symmetry is no longer restored after inflation and it is straightforward to check that

$$\rho_N/\rho_\phi \simeq N_0^2/\phi_0^2 \quad (22)$$

Although $\rho_N \sim \rho_\phi$ initially (we take $N_0 \sim \phi_0$), one can check that at the time of N decay the N oscillations dominate the energy density for $m_I > O(10^6)$ GeV. As in eq. (12) we write

$$\frac{R_{DN}}{R_R} \simeq m_I^{4/3} N_0^{2/3} / k^{4/3} \tilde{m}^{4/3} m_P^{2/3} \quad (23)$$

and from eq. (20), $n_B = k^4 \epsilon \phi_0^4 \tilde{m}^5 m_P^2 / m_I^4 m_G^2 N_0^2$ and s_D is again given by eq. (13) so that

$$\frac{n_B}{s} \propto k \epsilon \phi_0^4 \tilde{m}^{1/2} m_P^{1/2} / m_I m_G^2 N_0^2 \quad (24)$$

which for the same values of the parameters k , ϵ , etc., yields the limit $m_I > O(10^{12})$ GeV for $n_B/s > 10^{-11}$, as seen in the figure. If $m_I > 10^{14}$ (10^{16}) GeV and $T_R < 10^8$ - 10^9 GeV then $n_B/s \approx 10^{-14}$ (10^{-16}).

Other scenarios for low-energy baryosynthesis have been proposed, but we doubt that they can be realized in practice. One scenario [15] involves non-perturbative effects in the Weinberg-Salam model, but we believe that any such effects leading to $\Delta B \neq 0$ transitions are strongly suppressed at $T \gg m_W$ and $T \ll m_W$, and see no reason why they should be important when $T \approx m_W$ [16]. Another scenario [17], involving the late out-of-equilibrium decays of ν^c or N fields, requires many *ad hoc* assumptions about their masses and couplings which do not seem to be realized in specific models [8, 9].

Another question which must be addressed is how the intermediate-scale breaking is to be achieved. Normally [7] in supergravity theories one drives some supersymmetry-breaking scalar (mass)² $\tilde{m}_N^2$ negative as in eq. (6) using radiative corrections resummed using the matrix of renormalization group equations:

$$\mu \frac{\partial \tilde{m}^2}{\partial \mu} = O\left(\frac{\alpha}{\pi}\right) \tilde{m}^2 \quad (25)$$

over the large energy range $m_P > \mu > \tilde{m} \approx 10^3$ GeV. The magnitude of the corresponding vacuum expectation value, like $\langle 0|H|0 \rangle$ in the Standard Model, is then $\lesssim \mu_0$, where μ_0 is fixed by the condition

$$m_N^2(\mu_0) = 0: \quad \mu_0/m_P = \exp\left[-O(1)/\alpha\right] \quad (26)$$

This scenario cannot be applied to intermediate-scale models because the field N has no renormalizable interactions at scales below m_I . The self-couplings of the N field are by assumption non-renormalizable as in eq. (6). Any renormalizable superpotential coupling of the N field to other fields, e.g. NDD^c , gives large masses $O(m_I)$ to the D/D^c particles, which therefore cannot contribute to the renormalization group equations at scales below $O(m_I)$. Non-renormalizable interactions can only contribute to $\mu \partial \tilde{m}^2 / \partial \mu$ in order $\tilde{m}^{2+n}/m_P^n$, and it is easy to convince oneself that they cannot drive $\tilde{m}^2 < 0$ at a scale $\mu \gtrsim m_I$. Therefore, the evolution of $\tilde{m}_N^2$ effectively ceases at $\mu = O(m_I)$, and it must be that $\mu_0 \gtrsim m_I$. But we argued previously that m_I must be $\gtrsim 10^{16}$ GeV, so the exponential hierarchy in eq. (26) is unlikely to be relevant. Indeed, the logarithmic range between m_P and m_I is

so small that renormalization group resummation of the radiative corrections may not be appropriate. Is there any way of rescuing the usual scenario of symmetry breaking by radiative corrections? Resolution of this issue requires a deeper understanding of supersymmetry breaking in the observable sector.

It should be noted that a typical intermediate scale has not one but many flat directions, and if any of the corresponding scalar fields can acquire a negative mass-squared as in eq. (6), many might do so. In this case, domains of the Universe which were causally separated when the temperature T fell below $O(\tilde{m})$ would choose different vacua. It is easy to convince oneself that the probability for tunnelling from a false vacuum to the lowest of all these minima is negligible, so the Universe gets stuck with an unacceptable domain structure. One might have hoped that the choice of minimum could take place before or during inflation, so that the entire visible Universe could sit in one of the local minima. However, we believe that this scenario is not possible because of the de Sitter fluctuations [14] mentioned earlier, which send the flat fields to random values which are $O(m_{\text{Pl}})$ from any minimum.

We conclude with some comments on baryon and lepton number violation and phenomenological problems such as $H_u H_d$ and $H_u L$ mixing in the effective low-energy theory. Using a standard notation for the fields in a 27 representation of E_6 , the possible renormalizable trilinear terms in the superpotential are

$$P \supset \left\{ \begin{array}{l} q d^c \bar{H}, \quad q u^c H, \quad l e^c \bar{H}, \quad l \nu^c H, \quad H \bar{H} N, \quad D D^c N, \\ D q q, \quad D^c u^c d^c, \quad D^c l q, \quad D e^c u^c, \quad D r^c d^c \end{array} \right\} \quad (27)$$

Each of the couplings λ_i in eq. (27) should be understood as a matrix in generation space, with the range of indices extended to include light survivors from split multiplets, if they exist. There is a maximal $SU(6) \times SU(2)_N$ subgroup of E_6 with respect to which the (H_u, q, u^c, e^c) are $SU(2)_N$ singlets, and the $(H_d, L), (\nu^c, N)$, and (D, d^c) are $SU(2)_N$ doublets. If only one of $\langle 0 | \bar{\nu}^c, N | 0 \rangle \neq 0$, which we may assume without loss of generality to be $\langle 0 | N | 0 \rangle \neq 0$, there are no masses for the L and d^c , which we interpret as lepton doublets and charge $+1/3$ antiquarks, respectively. Phenomenology also requires a pair of light Higgs doublets H_u and H_d , which can only occur if there are the appropriate zeros in the coupling matrix λ_6 of eq. (27). Even if there are such zeros, generating an $H_u H_d$ mixing coefficient of the appropriate magnitude, $O(m_w)$ is not trivial [9]. Specifically, (at least) one pair, H_u and H_d , must not be allowed to couple to the N which receives the intermediate v.e.v., and either the model must have the possibility of an additional v.e.v. for a light N' with $\langle 0 | N' | 0 \rangle \approx 1 \text{ TeV}$ which produces the necessary $H_u H_d$ mixing, or we must rely on the presence of a non-renormalizable coupling such that the $H_u H_d$ mixing is of the order of $m_{\text{Pl}}^{p+1}/m_{\text{Pl}}^p$. It should be noted, however, that if there are zeros in the λ_6 matrix, there is also the possibility of zeros in the λ_8 matrix of eq. (27), resulting in light D/D^c quarks, which then could violate baryon number through the couplings $\lambda_4, \lambda_9, \lambda_{10}$, and λ_{11} . If both $\langle 0 | \bar{\nu}^c, N | 0 \rangle \neq 0$, then in general the light spectrum contains, as light antiquarks, arbitrary mixtures of d^c and D^c , the light leptons are mixtures of L and

H_d , and the light $Y = -1/2$ Higgses are orthogonal mixtures of H_d and L [9]. If one can overcome the problem of generating $H_u H_d$ mixing described above, there is also the likelihood of generating $H_u L$ mixing. As in $H_u H_d$ mixing discussed above, an intermediate-scale v.e.v. for $\bar{\nu}^c$ will be likely to produce $H_u L$ mixing of similar order, although phenomenological limits on such mixing require it to be $\leq O(1)$ GeV. Therefore, the effective low-energy theory is likely to violate maximally both B and L conservation^{*)}. Reference [9] contains a study of the effective low-energy theories obtained from different low-energy limits of the Yau models [8], corresponding to large intermediate-scale symmetry breaking in different flat directions. We find that such models with both $\langle 0|\bar{\nu}^c, N|0\rangle \neq 0$ do indeed normally violate B and L maximally, although this can be avoided by a careful [8] choice of flat direction which lacks *a priori* motivation.

We are not so foolish as to claim a no-go theorem for intermediate-scale models, but we believe that some studies in the literature have taken a Panglossian attitude which understates their deep problems.

Acknowledgements

The work of K.E. and 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.

^{*)} The latter is hardly surprising if $\langle 0|\bar{\nu}^c|0\rangle \neq 0$.

REFERENCES

- [1] P. Candelas, G. Horowitz, A. Strominger and E. Witten, Nucl. Phys. **B258** (1985) 46;
E. Witten, Nucl. Phys. **B258** (1985) 75.
- [2] M. Dine, V. Kaplunovsky, M. Mangano, C. Nappi and N. Seiberg, Nucl. Phys. **B259** (1986) 519;
S. Cecotti, J.P. Derendinger, S. Ferrara, L. Girardello and M. Roncadelli, Phys. Lett. **156B** (1985) 318;
J.D. Breit, B.A. Ovrut and G. Segré, Phys. Lett. **158B** (1985) 33.
- [3] Y. Totsuka, Proc. Int. Symposium on Lepton and Photon Interactions at High Energies, Kyoto, 1985, eds. M. Konuma and K. Takahashi (Kyoto Univ, Kyoto, 1985), p. 120.
- [4] B. Campbell, J. Ellis and D.V. Nanopoulos, Phys. Lett. **141B** (1984) 229.
- [5] I. Affleck and M. Dine, Nucl. Phys. **B249** (1985) 361.
- [6] B. Campbell, J. Ellis, D.V. Nanopoulos and K. Olive, Mod. Phys. Lett. **A1** (1986) 389;
J. Ellis, D.V. Nanopoulos and K. Olive, CERN preprint TH. 4552 (1986).
- [7] For a recent review, see: A.B. Lahanas and D.V. Nanopoulos, Phys. Reports **145** (1987) 1.
- [8] B.R. Green, K.H. Kirklin, P. Miron and G.G. Ross, Phys. Lett. **180B** (1986) 69; Nucl. Phys. **B278** (1986) 667 and paper in preparation.
- [9] J. Ellis, K. Enqvist, D.V. Nanopoulos and K. Olive, paper in preparation.
- [10] S. Weinberg, Phys. Rev. **D26** (1982) 287;
N. Sakai and T. Yanagida, Nucl. Phys. **B197** (1982) 533.
- [11] K. Yamamoto, Phys. Lett. **168B** (1986) 341.
- [12] K. Enqvist, D.N. Nanopoulos and M. Quiros, Phys. Lett. **169B** (1986) 343.
- [13] O. Bertolami and G.G. Ross, Oxford University Department of Theoretical Physics preprint 70/86 (1986).
- [14] T.S. Bunch and P.C.W. Davies, Proc. Roy. Astr. Soc. **360** (1978) 117;
A. Vilenkin and L.A. Ford, Phys. Rev. **D26** (1982) 1231;
A.D. Linde, Phys. Lett. **116B** (1982) 335;
A.A. Starobinsky, Phys. Lett. **117B** (1982) 175.
- [15] V.A. Kuzmin, V.A. Rubakov and M.E. Shaposhnikov, Phys. Lett. **155B** (1985) 83;
M.E. Shaposhnikov, Inst. for Nucl. Research, Moscow preprint (1986).
- [16] K. Enqvist, D.V. Nanopoulos and S. Rao, Univ. of Wisconsin-Madison preprint MAD/TH/86-32 (1986);
J. Ellis, R.A. Flores, S. Rudaz and D. Seckel, paper in preparation.
- [17] K. Yamamoto, Johns Hopkins University preprint (1986);
R.N. Mohapatra and J. Valle, University of Maryland preprint (1986).

The domain of the (m_I, T_R) plane which is allowed by the constraint $n_B/s > 10^{-11}$ in the Affleck-Dine scenario for baryosynthesis.