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SUPER B PRODUCTION, MATTER STABILITY AND GAUGE HIERARCHY
IN SUSY GUTS

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ABSTRACT

We analyze the connections between the problems of the cosmological baryon asymmetry (CBA), matter stability and mass hierarchy in supersymmetric Grand Unified models. We show that the typical delay of the phase transition in supersymmetric theories as well as a natural solution to the triplet-doublet Higgs hierarchy problem imply baryon production after the Grand Unification phase transition which, in a wide class of SUSY GUTs takes place at $T \sim 10^9 - 10^{10}$ GeV. Light Higgs colour triplets as a means of producing the CBA are discussed in detail. An alternative mechanism involving a singlet superfield is proposed leading to striking consequences in proton decay with the appearance of (B+L) conserving decay modes with a muon and a kaon in the final state: $n \rightarrow \mu^- K^+$, $p \rightarrow \mu^- \pi^+ K^+$ and a lifetime of 10^{31} years.

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1. - INTRODUCTION

Globally supersymmetric theories¹⁾ possess unique renormalization properties²⁾ not shared by general renormalizable field theories. As has recently been vigorously emphasized, the absence of quadratic divergences as well as powerful non-renormalization theorems lead to the solution of the technical side of the gauge hierarchy problem³⁾, namely the stability of scalar masses under radiative corrections. Thus, in supersymmetric Grand Unified Theories⁴⁾, any fine adjustment of parameters made at the tree level will be respected by radiative corrections to all orders in perturbation theory. More specifically, in an $SU(5)$ Grand Unified Theory, the $SU(2)$ Higgs doublets, whose expectation value will induce the electroweak breaking, must have a negative mass squared of $O(100 \text{ GeV})^2$, while their colour-triplet partners must have a superheavy positive mass of at least 10^{10} GeV in order to avoid too fast a proton decay. The masslessness of Higgs doublets (100 GeV is essentially zero compared with 10^{10} GeV) is usually achieved through fine tuning of the parameters at the tree level and then supersymmetry is invoked to keep the Higgs mass hierarchy stable to all orders in perturbation theory. This is, however, highly unnatural. As an alternative, leading to naturally massless Higgs doublets, we recently proposed⁵⁾ a non-minimal supersymmetric $SU(5)$ model making use of a superheavy supermultiplet in the $\underline{50}$ representation, which contains only coloured triplets, but no iso-doublets. Allowing the usual triplets contained in the fundamental $\underline{5}$ and $\bar{\underline{5}}$ to obtain mass only through their couplings to $\underline{50}$ (and $\bar{\underline{50}}$), we end up with heavy triplets and massless doublets without any fine tuning. It is remarkable, however, that the diagonalization of the triplet mass matrix can lead us naturally to colour triplets of an intermediate mass, roughly 10^{10} GeV . It is no coincidence that such "light" triplets can provide us with a mechanism for the generation of the cosmological baryon asymmetry.

Grand Unified Theories describe the evolution of the early Universe⁶⁾ in terms of successive phase transitions corresponding to the breakdowns suffered by the unifying gauge group. One of the outstanding successes of GUTs is the dynamical generation of the cosmological baryon asymmetry (CBA), i.e., the universal asymmetry between matter and antimatter. Supersymmetry drastically affects the picture of the early Universe. Supersymmetric GUTs lead to cosmological scenarios quite different from those encountered in ordinary GUTs. As a general feature, the phase transition corresponding to the breakdown of the grand unifying group is postponed until temperatures much lower than the unification energy scale. In the cosmological scenario⁷⁾⁻⁹⁾ recently proposed by two of us (K.T. and D.V.N.) and independently by M. Srednicki, the phase transition can take place at temperatures as low as $10^9 - 10^{10} \text{ GeV}$. This has a dramatic influence on the possible mechanisms for the creation of the CBA¹⁰⁾. In particular, the generation

of the CBA through coloured Higgs boson decays requires "light" Higgs triplets of mass (roughly) 10^{10} GeV. Quite generally, an "overdue" phase transition requires the existence of intermediate mass particles for the dynamical generation of CBA.

The increase in the number of fundamental particles that accompanies supersymmetrization leads to an increase in the unification mass scale by one or two orders of magnitude [this is not without an exception however¹¹⁾]. Therefore, nucleon decay through gauge boson exchange is suppressed. Nevertheless, Higgs mediated proton decay occurs at a rate compatible with present experimental limits¹²⁾. The decay modes are characteristic of both supersymmetry and the particular features of the Higgs sector. As a result, there is an intimate connection between nucleon decay modes and the type of processes that can create a CBA.

The paper is organized as follows. In Section 2 we discuss phase transitions in supersymmetric GUTs and their implications on the creation of the CBA. In Section 3 we concentrate on "light" coloured Higgs triplets as a means to generate the CBA. In Section 4 we discuss matter stability in connection with the coloured Higgses employed for the production of the baryon asymmetry. In Section 5 we propose an alternative mechanism for the generation of the CBA through a supersinglet. In Section 6 we examine a SUSY GUT version of the supersinglet mechanism with important implications on proton decay. Finally, in Section 7, we state our conclusions.

2. - PHASE TRANSITIONS AND CBA IN SUPERSYMMETRIC GUTS

The phase transitions from the grand unifying phase to the phase with broken gauge symmetry is generally more problematic in supersymmetric GUTs than in standard GUTs. As shown by two of us (D.V.N. and K.T)⁷⁾, the minimal $SU(5)$, where only a $\underline{24}$ Higgs supermultiplet is involved in the $SU(5)$ symmetry breaking, presents a perverse reluctance to abandon the symmetric phase. It may be useful to remind the reader of the major reasons that led to this conclusion. At $T = 0$, the superpotential for the $\underline{24}$ supermultiplet Φ is

$$W = \frac{\lambda}{3} \text{Tr} [\Phi^3] + \frac{M}{2} \text{Tr} [\Phi^2] \quad (1)$$

leading to the potential

$$V = \text{Tr} \left| M \Phi + \lambda \left(\Phi^2 - \frac{1}{5} \text{Tr} (\Phi^2) \right) \right|^2 + \frac{e^2}{2} \text{Tr} ([\Phi, \Phi^\dagger])^2 \quad (2)$$

V vanishes for the minima

$$\begin{aligned}
 \langle \Phi \rangle &= 0 && [\text{SU}(5) \text{ phase}] \\
 \langle \Phi \rangle &\propto \text{Diag}(1, 1, 1, 1, -4) && [\text{SU}(4) \times \text{U}(1) \text{ phase}] \\
 \langle \Phi \rangle &\propto \text{Diag}(2, 2, 2, -3, -3) && [\text{SU}(3) \times \text{SU}(2) \times \text{U}(1) \text{ phase}]
 \end{aligned} \tag{3}$$

When the temperature T is turned on, the degeneracy is removed, since supersymmetry is manifestly violated at finite temperature¹³⁾, but unfortunately in the "wrong" way. For temperatures above the scale of supersymmetry breaking, the global minimum is the $\text{SU}(5)$ symmetric one, while the other two phases correspond only to local minima. Indeed, at high temperatures, the potential becomes

$$V(\phi, T) \simeq V(\phi) + c_1 T^2 + c_2 T^4 \tag{4}$$

where

$$c_2 = -\frac{\pi^2}{90} (N_B + \frac{7}{8} N_F) \tag{5a}$$

and

$$c_1 = \frac{1}{24} T_r \{ m_S^2 + m_F^2 + 3 m_V^2 \} \tag{5b}$$

In Eq. (5a) we have denoted by N_B and N_F the number of helicity states of light bosons and fermions, respectively. A calculation of the $\text{SU}(5)$ β function shows that the $\text{SU}(5)$ gauge interactions become strong at $\Lambda_{\text{SU}(5)} \sim O(10^9 - 10^{10})$ GeV. The presence of strong gauge forces undoes the naive scenario of an exceptionally late $\text{SU}(5)$ phase transition^{7),8)}. Although we cannot make any quantitative predictions as to the strong coupling regime, it appears very reasonable that the confining forces will, as a result, cause a decrease in the number of light $\text{SU}(5)$ degrees of freedom. At the same temperature on the other hand, the $\text{SU}(3) \times \text{SU}(2) \times \text{U}(1)$, as well as the $\text{SU}(4) \times \text{U}(1)$ phase, is still weakly coupled. The coefficient c_2 of (5a) will decrease for $\text{SU}(5)$, thus making the broken phases preferable. It is interesting to note that although in the minimal case both broken phases have the same free energy, in the case of four pentaplets of Higgses (the lowest number capable of generating a baryon asymmetry via colour triplets) the $\text{SU}(3) \times \text{SU}(2) \times \text{U}(1)$ case is preferred⁹⁾. Nevertheless, it has been argued that although the phase transition to the broken phases can proceed, it might take

a very long time, since the barrier separating the phases is in general much bigger than the range of strong coupling phenomena⁸⁾. This can be bypassed by complicating the model so that paths exist in field space with a negligible barrier⁹⁾. A more elaborate choice of scalar supermultiplets including SU(5) singlet superfields could achieve it. Hence, quite generally, it can be said that in contrast with what happens in usual GUTs, we expect a large class of supersymmetric GUTs to display a typical delay in the phase transition. Our analysis of the CBA in supersymmetric GUTs concerns this wide class of models in which the transition occurs many orders of magnitude below M_X . Of course, it is possible to arrange it so that the Universe is in the broken phase already at M_X . For this class of models nothing much new can be said about CBA apart from what is already known from standard GUTs.

The delay in the phase transition should not be confused with any supercooling phenomenon since, as long as supersymmetry is unbroken, the vacuum energy is purely thermal and no exponential expansion is possible. In addition, no substantial reheating is expected when the phase transition is over. Some reheating is present due to the difference in degrees of freedom before and after, but it is not relevant. The question of the presence of any reheating is a central issue in the problem of CBA. In the case of supercooled ordinary GUTs, the CBA can be generated only after the accomplishment of the phase transition since the reheating widely dilutes any pre-existing CBA. Thus, in supersymmetric GUTs, in spite of the delay of the phase transition to $10^9 - 10^{10}$ GeV, both the following two possibilities for the generation of the CBA are viable in principle:

- a) CBA before the phase transition
- b) CBA during or after the phase transition.

Clearly, the first option requires the presence of non-thermalized modes¹⁴⁾ with a net ΔB . As is well known, this can be the case of SU(5), where fermions are arranged to the reducible representation $\underline{\bar{5}} + \underline{10}$. Gauge interactions conserve the number of fermionic 5-plets and 10-plets ("fiveness" and "ten-ness") separately. This means that any "net number" of 5-plets and 10-plets cannot be erased by gauge boson-mediated SU(5) interactions. However, Higgs boson couplings to fermions do not conserve "fiveness" and "ten-ness" separately, but only a global U(1) symmetry which is a linear combination of those two symmetries. So, unless the Higgs coloured triplets in SU(5) get a mass and decay well before the SU(5) phase transition, it is not possible to generate a net ΔB before such a transition, because the Higgs mediated interactions spoil the exact conservation of "fiveness" and "ten-ness". At first sight, we would say that both options for the mass of Higgs coloured triplets, i.e., that they get mass either before or after the SU(5)

breaking, are available and no general principle favours one choice over the other. As we shall argue in the next Section, the search for a natural solution to the triplet-doublet problem⁵⁾ (Higgs hierarchy problem) points in favour of a triplet mass only after the SU(5) breaking. In such a case, we are forced to generate the CBA only after the phase transition. Thus, since the phase transition will be typically delayed, we will have the creation of CBA at an intermediate scale.

3. - THE CASE OF "LIGHT" COLOUR TRIPLETS IN SU(5)

Let us first remind ourselves of the "triplet-doublet" problem in SU(5). The relevant terms in the superpotential are

$$\mathcal{W} \simeq \alpha H_{\bar{5}} \Phi_{24} H_5 + m H_{\bar{5}} H_5 + \dots \quad (6)$$

Along the $SU(3) \times SU(2) \times U(1)$ direction, we have

$$\langle \Phi_{24} \rangle = \frac{m'}{3a} \text{Diag}(2, 2, 2, -3, -3) \quad (7)$$

Then,

$$\mathcal{W} = (m + \frac{2}{3}m') H_{\bar{3}} H_3 + (m - m') H_{\bar{2}} H_2 \quad (8)$$

In order to keep the Higgs doublets massless, we must adjust m to be exactly equal to m' . Technically, there is no problem since the non-renormalization theorem of the superpotential ensures that any tree level adjustment will be respected by radiative corrections to all orders in perturbation theory. Nevertheless, such an incredibly accurate adjustment of two otherwise arbitrary parameters appears highly unnatural.

Before searching for a solution to the triplet-doublet problem, we are faced with two logical possibilities:

- a) either the Higgs 5-plets possess a direct mass term in the initial potential and, after the SU(5) breaking, some compensation occurs for the doublet masses so that they become negligible in comparison to the triplet masses, or
- b) the 5-plets are massless before the SU(5) breaking and, for some reason, after the SU(5) breaking, only the triplets get mass while the doublets remain massless^{*)}.

*) In analogy with the electron and neutrino mass in the standard model.

We have recently proposed a solution along the lines of strategy b)⁵⁾. Since the question is of particular interest for the problem of CBA, we briefly recall the major points of our solution. The crucial point is to avoid the introduction of a direct mass term for the $\underline{5}$ and $\bar{\underline{5}}$ and of representations that contain colour-singlet, iso-doublet components. On the contrary, we allow direct mass terms for representations that contain colour triplets but no iso-doublets. This is the case of the representation $\underline{50}$ of $SU(5)$ which decomposes under $SU(3) \times SU(2) \times U(1)$ as

$$\underline{50} = (\underline{8}, 2) + (\underline{6}, 3) + (\bar{\underline{6}}, 1) + (\underline{3}, \bar{2}) + (\bar{\underline{3}}, 1) + (1, 1) \quad (9)$$

$\underline{50}$ [and $\bar{\underline{50}}$ for an anomaly-free supersymmetric $SU(5)$] mix with $\underline{5}$ and $\bar{\underline{5}}$ through trilinear couplings involving a Higgs supermultiplet in the $\underline{75}$ representation of $SU(5)$. $\underline{75}$ takes the place of the adjoint and its expectation value breaks $SU(5)$ uniquely to $SU(3) \times SU(2) \times U(1)$. The $\underline{50}$ and $\underline{75}$ multiplets were first introduced in ordinary $SU(5)$ in order to obtain a radiatively induced fermion spectrum¹⁵⁾. Denoting the $\underline{75}$ by Φ , the $\underline{50}$ by Θ and the $\underline{5}$ by H , we can write down a superpotential that reads

$$\begin{aligned} \mathcal{W} = & \frac{M}{2} \text{Tr}(\Phi^2) + \frac{a}{3} \text{Tr}(\Phi^3) + b \Theta \Phi H + c \bar{\Theta} \bar{\Phi} \bar{H} \\ & + \tilde{M} \bar{\Theta} \Theta + \dots \end{aligned} \quad (10)$$

After spontaneous symmetry breaking, the relevant mass terms are

$$\mathcal{W} \simeq \frac{bM}{a} \Theta_3 H_3 + \frac{cM}{a} \bar{\Theta}_3 \bar{H}_3 + \tilde{M} \bar{\Theta}_3 \Theta_3 + \dots \quad (11)$$

where

$$\langle \Phi \rangle \sim \frac{M}{a} \sim \mathcal{O}(10^{15} - 10^{16} \text{ GeV})$$

There is no mass term for $H_2, \bar{H}_2$ in (11) and thus the doublets remain massless. The mass matrix for the triplets is

$$\mathcal{M} = \begin{pmatrix} 0 & \tilde{m} \\ \tilde{m} & \tilde{M} \end{pmatrix} \quad (12)$$

where $\tilde{m} = bM/a \sim \mathcal{O}(10^{14} - 10^{15} \text{ GeV})$. As far as the value of $\tilde{M}$ is concerned, two choices are natural, namely $\tilde{M} = 10^{15} - 10^{16} \text{ GeV}$ or $\tilde{M} = M_{\text{Pl}} = 10^{19} \text{ GeV}$. The second option leads to the scenario envisaged in Ref. 12), where the mass eigenstates $H \simeq -(\tilde{m}/\tilde{M})\bar{\Theta}_3 + H_3$ have an intermediate mass roughly $\sim \mathcal{O}(10^{10} \text{ GeV})$. In

this case the Planck scale is responsible for the appearance of the intermediate scale 10^{10} GeV. It is remarkable that 10^{10} GeV is also the lower mass limit for coloured triplets imposed by proton decay¹⁶⁾. The other option leads to "usual" triplets of a mass 10^{15} GeV. As we shall see later, these two options lead to radically different scenarios for the origin of CBA.

In the previous section we reached the conclusion that a large class of supersymmetric GUTs display an "overdue" phase transition. On the other hand, our solution to the "triplet-doublet" problem requires no direct pentaplet mass before the breaking. This, as we explained, forces us to generate the CBA only after the phase transition. Since the transition is delayed down to scales of roughly $10^9 - 10^{10}$ GeV, it is clear that we must make the choice of the "light" triplets. Thus, the solution to the Higgs hierarchy problem together with a delayed phase transition point to the necessity of "light" triplets.

It is generally known that in any theory it is hard to generate a CBA at an intermediate mass scale, i.e., at a scale much below M_X . B violating interactions at a scale $\ll M_X$ can only erase any preceding CBA. They cannot produce a new CBA through decays because of the out-of-equilibrium condition. Apart from the fact that exceptions can be found to the above argument, it should be kept in mind that here we are dealing with a special type of phase transition. As we discussed in Section 2, the SU(5) symmetric Universe reaches the strong coupling regime at temperatures $\sim 10^9 - 10^{10}$ GeV. Then, SU(5) confinement sets in. In such a "phase" only SU(5) singlets are allowed. Thus, the coloured triplets form singlet bound states which subsequently decay⁹⁾. Assuming that the singlets live long enough, when they decay baryon violating scatterings and triplet annihilations will be ineffective. Hence, the triplets will decay out of equilibrium.

It is easy to convince oneself that a pair of pentaplets (a $\underline{5}$ and a $\bar{\underline{5}}$) are not enough to produce a non-zero CBA¹⁰⁾. Similarly to what happens in the ordinary SU(5) model with only a 5-plet, the imaginary part of the graph shown in Fig. 1

$$Im Tr(abc d)$$

would vanish since $a = d^+$ and $b = c^+$. We must add at least a second pair of 5-plets with different couplings to quarks and leptons. The mass of the second pair must not be too different since the generated baryon asymmetry is suppressed proportionally to the mass ratio squared. If all four triplets were to be super-heavy ($\sim M_X$), while all four doublets massless, $\sin^2 \theta_W$ would come out

unacceptably large¹⁷⁾. It is always possible, however, to give a full pentaplet a superheavy mass in which case its associated iso-doublet does not contribute to the mixing angle¹⁸⁾. We certainly do not have to resort to this trick since for semi-superheavy triplets ($\sim 10^{10}$ GeV) $\sin^2\theta_W$ comes out in accordance with experiment even for four doublets¹²⁾, another good reason for "light" (10^{10} GeV) Higgs triplets.

4. - MATTER STABILITY AND THE CBA

The problem of matter stability (i.e., nucleon decay) is intimately connected with the mechanism one employs for the generation of CBA. As is well known by now, although gauge boson-mediated nucleon decay is unimportant, Higgs mediated-proton decay does occur at a sizeable rate. Apart from the usual dimension six operators that violate baryon number, supersymmetric theories predict proton decay via dimension five operators as well^{19),20),17)}. Graphs that give rise to dimension five operators are shown in Fig. 2. At first glance, coloured triplet Higgses of a mass of the order of 10^{10} GeV would lead to an unacceptably high proton decay rate if one takes the strength of supersymmetry breaking to be of the order of 100 GeV. Nevertheless, a moderate increase in the supersymmetry breaking scale [$\sim (10-100)$ TeV] can prevent too fast a proton decay rate¹⁰⁾. In particular, the limits obtained by E.N.R.¹⁷⁾ are

$$\frac{\mu_{sq}^2 \mu_H}{\mu_{\tilde{W}}} \gtrsim (0.4-4) 10^{18} \text{ GeV}^2 \quad (13a)$$

$$\frac{\mu_{sq}^2 \mu_H}{\mu_{\tilde{g}}} \gtrsim (4-40) 10^{18} \text{ GeV}^2 \quad (13b)$$

where m_{sq} stands for the squark mass and $m_{\tilde{g}}, m_{\tilde{W}}$ stand for the wino and gluino Majorana masses, respectively. Then, for $m_H \sim 10^{10}$ GeV, we get

$$\mu_{sq}^2 / \mu_{\tilde{W}} \gtrsim (0.4-4) \cdot 10^8 \text{ GeV} \quad (14a)$$

$$\mu_{sq}^2 / \mu_{\tilde{g}} \gtrsim (4-40) \cdot 10^8 \text{ GeV} \quad (14b)$$

For $\tilde{W}$ and $\tilde{g}$ Majorana masses in the 10 GeV range, we see that $m_{sq} \sim (10-100)$ TeV is sufficient to match the bounds (14). So, dimension five operators are no longer dangerous. On the other hand, the usual dimension six operators¹²⁾ due to Higgs exchange (Fig. 3) lead to a rate compatible with the experimental limit as long as $m_H \gtrsim 10^{10}$ GeV as is known from the "old" analysis of EGN¹⁶⁾.

Nevertheless, raising the supersymmetry breaking scale does not seem to be the most appealing way to render dimension five operators innocuous. As we stressed in the previous chapter, at least two pairs of Higgs supermultiplets are needed in order to generate a non-vanishing CBA via Higgs triplets. If one of them couples exclusively to the 1 generation (which does not seem an artificial requirement in view of the fact that the third generation fermions have distinctively higher masses than those of the first and second generations) then, dimension five operators, although present, can be energetically forbidden for the decay of ordinary nucleons¹⁰⁾. In particular, let us introduce two pairs of Higgs 5-plets $H^{(1)}, \bar{H}^{(1)}$ and $H^{(2)}, \bar{H}^{(2)}$. We restrict the Yukawa couplings as follows

$$\begin{aligned} \mathcal{W} \sim & d_{ij} Q_{10}^{(i)} Q_{10}^{(j)} H^{(1)} + d' Q_{10}^{(3)} Q_{10}^{(3)} H^{(2)} \\ & + f_{ij} Q_{10}^{(i)} Q_{\bar{5}}^{(j)} \bar{H}^{(2)} + f Q_{10}^{(3)} Q_{\bar{5}}^{(3)} \bar{H}^{(1)} \end{aligned} \quad (15a)$$

Furthermore, the mass terms for the triplets are

$$m_1 \bar{H}_3^{(1)} H_3^{(1)} + m_2 \bar{H}_3^{(2)} H_3^{(2)} \quad (15b)$$

A mixing term $\bar{H}^{(1)} H^{(2)}$ or $\bar{H}^{(2)} H^{(1)}$ is not present. It can be generated only after supersymmetry breaking and it will necessarily be of $O(m_s)$. Now, reconsider the diagrams of Fig. 2. We shall draw our attention only to the case in which the gauge fermion λ is a gluino or a photino, since only these particles remain massless if the supersymmetry breaking does not include a Majorana mass for them [the winos and binos get a Dirac mass in combination with Higgs fermions during the SU(2) super-Higgs effect, so that their Majorana mass can be zero]. An important point for what is to follow is that the gluino and photino exchange does not mix the generations^{*}). Because of the absence of $\bar{H}^{(1)} H^{(2)}$ or $\bar{H}^{(2)} H^{(1)}$ mixing terms, one of the down vertices of Fig. 2 must necessarily connect only $Q^{(3)}$ with $Q^{(3)}$, i.e.,

$$d Q_{10}^{(3)} Q_{10}^{(3)} H^{(2)} \quad \text{or} \quad f Q_{10}^{(3)} Q_{\bar{5}}^{(3)} \bar{H}^{(1)}$$

Since the most general type of mixing in SU(5), when all fermion masses are generated by 5's and $\bar{5}$'s, is

$$Q_{\bar{5}} = \begin{pmatrix} \bar{d}^{(i)} \\ \bar{e}^{(i)} \\ \nu^{(i)} \end{pmatrix} \quad \text{and} \quad Q_{10} = \begin{pmatrix} (Uu)^{(i)} \\ e^{+(i)} & d^i & (\bar{u}U^+)^{(i)} \end{pmatrix} \quad (16)$$

(where U stands for the standard mixing matrix) then clearly one of the upper vertices will involve third generation fermions, something that renders proton decay kinematically impossible.

^{*}) This ensures the absence of flavour-changing neutral currents²¹⁾.

Of course, dimension six operators are present and for a Higgs mass roughly 10^{10} GeV they lead to proton decay (Fig. 3) at the usual rate of 10^{-31} years $^{-1}$ with the characteristic hierarchy of decay modes

$$\Gamma(\bar{\nu}_\mu k, \mu^+ k^0) : \Gamma(\bar{\nu}_e k, e^+ k^0) : \Gamma(e^+ \pi^0, \bar{\nu}_e \pi^+) \simeq 1 : \sin^2 \theta_c : (\sin^2 \theta_c)^2$$

where $\sin^2 \theta_c \simeq 1/20$ is the Cabibbo angle¹²⁾.

Let us now come to the generation of the CBA in the above scheme. Consider first a typical one-loop interference diagram (Fig. 4). Notice that only the diagonal elements d_{33} and f_{33} of the matrices f_{ij} and d_{ij} enter. This is clearly due to the restrictions we have imposed on the Yukawa Lagrangian (15) and to the basis (16) that has been chosen. Naively, one would expect the interference in Fig. 4 to be real. Indeed, if we can redefine the phases of d , f , d_{33} , f_{33} and of the mass terms $H^{(1)}\bar{H}^{(1)}$ and $H^{(2)}\bar{H}^{(2)}$, then the diagram has no chance of presenting an imaginary part. We have just six arbitrary phases in the definition of $Q_5^{(3)}$, $Q_{10}^{(3)}$, $H^{(1)}$, $\bar{H}^{(1)}$, $H^{(2)}$ and $\bar{H}^{(2)}$, so one could expect that any phase in Fig. 4 can be rotated away. On the contrary, these phases are connected to the phases that can give an imaginary part to the diagram of Fig. 4 through the specific combinations (15) and if one writes down the system of equations for the phases appearing in (15) one finds that not all the phases of the Yukawa couplings in Fig. 4 can be simultaneously removed. Thus the diagram gives rise to a CBA

$$\Delta B \propto \text{Im} (d_{33} f_{33}^* f d^*) \neq 0 \quad (17)$$

In the model of the previous section where a $\underline{50}$ -plet Θ was introduced to solve the doublet-triplet hierarchy problem, another diagram can be constructed for the generation of the CBA, where Θ plays a role⁵⁾. In order to get a non-vanishing ΔB we need two $\underline{50}$ supermultiplets. Their couplings to fermions are unrestricted

$$\mu_{ij} Q_{10}^{(i)} Q_{10}^{(j)} \bar{\Theta}^{(1)} + \mu'_{ij} Q_{10}^{(i)} Q_{10}^{(j)} \bar{\Theta}^{(2)} \quad (18)$$

and their masses come from the expectation value of the $\underline{75}$ supermultiplet Φ

$$b \Theta^{(1)} \Phi H^{(1)} + b' \Theta^{(2)} \Phi H^{(2)} + c \bar{\Theta}^{(1)} \Phi \bar{H}^{(1)} + c' \bar{\Theta}^{(2)} \Phi \bar{H}^{(2)} \quad (19)$$

as well as direct terms $\bar{\Theta}^{(1)}\Theta^{(1)}$, $\bar{\Theta}^{(2)}\Theta^{(2)}$. No terms $\bar{\Theta}^{(1)}\Theta^{(2)}$, $\bar{\Theta}^{(2)}\Theta^{(1)}$ appear. Then, the diagram of Fig. 5 presents a non-vanishing imaginary part

$$\text{Im Tr}(f_{i3}^* f'_{i3} h_{i3}) \neq 0 \quad (20)$$

and can give rise to a CBA.

One might wonder why we took such a strenuous effort to make the dangerous dimension five operators innocuous via a complicated Higgs system, when it is known that they can be avoided by imposing an extra symmetry, either an R symmetry²⁰⁾ or a gauged $\tilde{U}(1)$ symmetry¹⁹⁾. The reason is that, as we shall show shortly, the same global symmetry that forbids dimension five operators also forbids a non-zero CBA through the interference diagrams of the type of Fig. 1.

Independently of specific GUTs, the effective "low-energy" theory has to be a supersymmetric $SU(3) \times SU(2) \times U(1)$ invariant theory with the following chiral matter superfields

$$\begin{aligned} Q &= (3, 2, -1/6) \\ \bar{U} &= (\bar{3}, 1, -2/3) \\ \bar{D} &= (\bar{3}, 1, 1/3) \\ \bar{E} &= (1, 1, 1) \\ L &= (1, 2, -1/2) \end{aligned} \quad (21)$$

The operators of dimension five with $\Delta B \neq 0$ which can be built out of the superfields in (21) are

$$\begin{aligned} &\bar{D} D (Q_{i\alpha} Q_{j\beta} Q_{k\gamma} L_n) \epsilon_{\alpha\beta\gamma} \epsilon_{ijkn} \\ &\bar{D} D (Q_{i\alpha} Q_{j\beta} Q_{k\gamma} L_n) \epsilon_{\alpha\beta\gamma} (\tau^a \epsilon)_{ij} (\tau^a \epsilon)_{kn} \\ &\bar{D} D (\bar{D}_\alpha \bar{U}_\beta \bar{U}_\gamma \bar{E}) \epsilon_{\alpha\beta\gamma} \end{aligned} \quad (22)$$

where α, β, γ and i, j, k are $SU(3)$ and $SU(2)$ indices, respectively and $\bar{D} D$ is the usual projection operator whose dimension is one.

Next we introduce a global symmetry R with charge assignment

$$R(Q, \bar{U}, \bar{D}, \bar{E}, L) = -1$$

Then, all dimension five operators (necessarily of the type Φ_L^4) are forbidden. Automatically then one can see by inspection of Fig. 1 that, as long as R is unbroken, the one-loop Higgs triplet interference diagrams must vanish. The four internal fermion legs give rise to operators of the type (22) which vanish. Thus, we are led to the following theorem:

Any $\tilde{U}(1)$ (R or gauge) symmetry, that remains unbroken down to low energies, and under which quarks and leptons have $\tilde{U}(1)$ charges of the same sign, forbids the generation of CBA through colour triplet decays.

This holds true independently of any specific supersymmetric GUT.

In conclusion, if we impose an extra symmetry to kill the $\Delta B \neq 0$ operator of dimension five, then no CBA can originate from the decays of Higgs colour triplets. This point of course strengthens the importance of our Higgs system (15), but it can also motivate the search for alternative ways of generating CBA without employing colour triplet decays.

5. - AN ALTERNATIVE MECHANISM FOR THE GENERATION OF CBA

As is well known in ordinary GUTs, colour-triplet decays are not the only means of generating a CBA. In particular, a possible dominant role of superheavy fermion decays has been advocated by the authors in different contexts^{22),23)}. A supersymmetric version of the superheavy fermion scenario has been outlined in Ref. 7). Trying to be as general as possible, let us proceed by supposing that there exists a heavy (of intermediate mass $\sim 10^{10}$) $SU(3) \times SU(2) \times U(1)$ singlet superfield F . What is crucial for the role played by such a particle in the production of CBA is the way in which it can decay into ordinary light particles. We distinguish the following two possibilities:

- a) F cannot decay into light particles through effectively renormalizable decays, i.e., it can only decay through three-body channels with a process mediated by some superheavy particle (Fig. 6)²²⁾;
- b) F has at least one two-body channel to light particles (Fig. 7)²³⁾.

Option a) holds for those F 's whose combinations with ordinary fermions give rise to $SU(3) \times SU(2) \times U(1)$ quantum numbers which do not correspond to any of the light Higgs or gauge fields. Since, by assumption, the total decay rate for this class is dominated by three-body processes, ΔB is not suppressed by propagator

effects and it has been shown that a sizeable ΔB , in agreement with $N_B/N_Y \sim 10^{-9}$, can be obtained. In particular, it should be stressed that, depending on the ratio m_F/M_H , the F's can be prevented from decaying long after they have decoupled. In this case it is worth mentioning that for a certain period the Universe is "matter" dominated.

Let us now analyze in detail the second option. Here it is not too difficult to find an example of a heavy particle obeying the criterion stated in b). What has been termed in the literature as the right-handed neutrino, has an $SU(3) \times SU(2) \times U(1)$ invariant mass and presents a two-body decay channel to light particles given by

$$\nu_R \longrightarrow \nu_L + \phi^0 \quad (23)$$

where ϕ^0 is the neutral component of a light iso-doublet scalar. In analogy with the above, we introduce an $SU(3) \times SU(2) \times U(1)$ singlet superfield N . If the colour-triplet H_3 is lighter than N , then N presents two-decay channels

$$N \longrightarrow \ell + H_2, \quad N \longrightarrow q + H_3 \quad (24)$$

which carry a different B number. Notice that in this case we should keep in mind that i) the H_3 's must disappear very soon after the creation of the CBA in N decays since their B violating interactions in equilibrium can wash out the previous ΔB ; ii) the H_3 's cannot be lighter than the lower mass limit imposed by Higgs-mediated proton decay through dimension six operators ($\sim 10^{10} \text{ GeV}$). Although i) and ii) can be met, we find it more interesting to pursue an alternative type of N decay to those in (24).

Considering again an $SU(3) \times SU(2) \times U(1)$ effective theory, we can introduce another decay channel for N carrying $B \neq 0$

$$N \longrightarrow \bar{u} + \Phi \quad (25)$$

where Φ transforms as $(3, 1, 2/3)$ ^{*)}. Exercising our freedom over the couplings in (25), we can meet the out-of-equilibrium condition and finally obtain a sizeable ΔB . There is, however, an important consequence as far as proton decay is concerned. Φ cannot mediate proton decay directly as the usual colour triplets, but only through diagrams of the type in Fig. 8. This diagram is suppressed by

^{*)} Φ will eventually decay into a $\bar{d}d$ pair.

$(m_D^V/m_N)^2$ if we denote by m_D^V the Dirac entry in the neutrino mass matrix. Hence this process does not lead to an observable proton decay rate unless m_Φ is much less than m_N . For example, for $m_D^V/m_N \approx 10^{-10}$, we must have $m_\Phi \sim 10$ TeV in order to obtain a sizeable proton decay rate.

The introduction of Φ (and $\bar{\Phi}$ for no anomalies) suggests that, apart from the coupling $\bar{\Phi}\bar{D}\bar{D}$, which was proved to be not dangerous for proton decay, it would be interesting to ask whether a new Higgs supermultiplet could participate in a quark-lepton coupling. Indeed, a new Higgs superfield η , transforming as $(3, 2, 1/6)$ under $SU(3) \times SU(2) \times U(1)$, realizes the coupling²⁴⁾

$$L \bar{D} \eta \quad (26)$$

Notice again that the presence of η at an intermediate scale is also not dangerous for matter stability, since, analogously to what happens for Φ , η cannot mediate a direct proton decay. An $SU(3) \times SU(2) \times U(1)$ invariant coupling between η and Φ can be written down immediately

$$\epsilon_{\alpha\beta} (H_2)^\alpha \eta^\beta \bar{\Phi}_i \quad (27)$$

where H_2 is the usual Higgs iso-doublet.

It is readily obvious from (27) that B-L is violated in this model. The B-L content of $\bar{\Phi}$ and η is $2/3$ and $4/3$ respectively. On the other hand, the combination B+L is conserved. Then proton decay through the coupling (27) would violate B-L and conserve B+L. One can easily find a proton decay diagram (Fig. 9) involving (27). Such a process predicts a lepton in the final state (instead of an antilepton as in B-L conserving theories). J.F. Nieves and two of the authors (A.M. and T.Y.) have recently analyzed in detail the consequences of proton decay through these processes. Keeping Φ and η at some intermediate scale below 10^{10} GeV it is possible to have the B+L conserving process as the dominant decay mode.

6. - (B-L) VIOLATION FROM SUPERSYMMETRIC SU(5)

The striking signature of B+L conservation in nucleon decay and thus of the presence of a lepton instead of an antilepton in the final state, is motivation enough (apart from its relevance in the origin of CBA) to construct a supersymmetric GUT with this property. The $SU(3) \times SU(2) \times U(1)$ assignments of Φ and $\bar{\eta}$ ($\bar{\Phi}$ and η) hint that they are members of the same supermultiplet of a grand

unifying group. In fact they are contained in the $\underline{10}$ (and $\overline{10}$) supermultiplet of $SU(5)$. The couplings allowed by $SU(5)$ gauge symmetry are

$$\mathcal{W} \sim f_{ij} Q_5^{(i)} Q_5^{(j)} \chi_{10} + \mu \chi_{10} \chi_{10} + h \chi_{10} \chi_{10} H_5 + \dots \quad (28)$$

where i, j are generation indices.

$SU(5)$ singlets can be easily incorporated with couplings

$$N Q_{10} \chi_{10} + \dots \quad (29)$$

Statistics compels us to choose f_{ij} in (28) as antisymmetric. Thus, the Higgs-matter vertex in (28) will involve the second generation and the lepton in the final state will be a muon. Similarly, the $\overline{D}\overline{D}$ piece will involve strange quarks. Therefore we also expect a kaon in the final state. Summarizing, this variety of $SU(5)$ leads to nucleon decay with a muon and a kaon in the final state. Depending on the parameters ($\mu < 10^{10}$ GeV for CBA), the "exotic" modes

$$p \rightarrow \mu^- \pi^+ k^+ , \quad n \rightarrow \mu^- k^+$$

could dominate.

Independently of nucleon decay, there is another striking property of the $\underline{10}$ that should be stressed. As is the case for any representation, if all its components are at the same intermediate mass, then its effect on the $SU(3)$, $SU(2)$ and $U(1)$ β functions is such that the combinations entering in $\sin^2 \theta_W$ and M_X remain unchanged. This means that our model with a $\underline{10}$ -plet whose components are all at a mass $< 10^{10}$ GeV does not depart from the successful prediction of the minimal $SU(5)$ *).

In conclusion, we would like to stress the important points of the last two sections.

- i) The decay of the supersinglet N can give rise to the desired CBA at temperatures below $M_X (\sim 10^9 - 10^{10}$ GeV), without any problem for the out-of-equilibrium condition.

*) Notice that if only the iso- and colour-singlet component of $\underline{10}$ remains massless while all the other components are at M_X , then only the $\widetilde{U}(1)$ β function is modified in such a peculiar way that one recovers exactly the same expressions for M_X and $\sin^2 \theta_W$ as in ordinary $SU(5)$. This was noticed independently by the two groups of Ref. 11).

- ii) B violation is mediated to the supersinglet via a superfield χ_{10} with an inferior mass. Since χ_{10} does not mediate direct proton decay, its presence at a low scale does not create any fear for matter stability.
- iii) Observable proton decay can arise due to χ_{10} dominated by the (B+L) conserving modes

$$p \rightarrow \pi^+ k^+ \mu^- \quad n \rightarrow \mu k^+$$

7. - CONCLUSIONS

Our major points on the connections between the creation of the CBA, the conditions on matter stability and the hints coming from the Higgs hierarchy problem in supersymmetric Grand Unified models, can be summarized as follows.

- 1) In most SUSY GUTs, the symmetric phase persists down to temperatures much lower than the Grand Unification scale. In the case of minimal SU(5), the Universe is still in the SU(5) symmetric phase down to temperatures of order $10^9 - 10^{10}$ GeV where $\alpha_G(T) \sim O(1)$. Thus, we are led to pursue the creation of the CBA below an intermediate mass scale around $10^9 - 10^{10}$ GeV.
- 2) In a supersymmetric theory, the vacuum energy density is purely thermal. Therefore, we do not have any period of exponential expansion and do not expect any appreciable reheating. Thus, in principle, one could expect that both the possibilities of producing the CBA before or after the phase transition are available. Nevertheless, a natural solution of the triplet-doublet mass hierarchy problem favours the generation of the CBA via colour-triplet decays after the phase transition. This comes about because the only way known of obtaining massless Higgs doublets without any fine tuning requires that the mass of the associated triplets be zero also before the breaking.
- 3) The use of colour-triplets of a low mass ($\sim 10^{10}$ GeV) for the creation of the CBA requires the suppression of dimension five operators in nucleon decay. This can be achieved only with the choice of a special Higgs system. Dimension five operators cannot be forbidden with the imposition of some $\tilde{U}(1)$ symmetry since the same symmetry kills the Higgs interference diagrams that give a CBA.

- 4) Alternatives to colour-triplet decays exist. We propose a scenario for the creation of the CBA in terms of a singlet superfield of an intermediate mass ($\sim 10^{10}$ GeV) which can decay into coloured particles of inferior mass ($< 10^{10}$ GeV). These particles pose no fear for matter stability since they cannot mediate direct proton decay. Nevertheless, they can have (B-L) violating interactions which can lead to nucleon decay within the experimental limit dominated by exotic (B+L)-conserving modes involving a muon and a kaon in the final state.

REFERENCES

- 1) Y.A. Gol'fand and E.P. Likhtman, JETP Lett. 13 (1971) 323;
D. Volkov and V.P. Akulov, Phys. Lett. 46B (1973) 109;
J. Wess and B. Zumino, Nucl. Phys. B70 (1974) 39;
For a review, see:
P. Fayet and S. Ferrara, Phys. Rep. 32 No. 5 (1977) 249.
- 2) J. Iliopoulos and B. Zumino, Nucl. Phys. B76 (1974) 310;
S. Ferrara, J. Iliopoulos and B. Zumino, Nucl. Phys. B77 (1974) 413;
S. Ferrara and O. Piguet, Nucl. Phys. B93 (1975) 261.
- 3) E. Gildener and S. Weinberg, Phys. Rev. D13 (1976) 3333;
A.J. Buras, J. Ellis, M.K. Gaillard and D.V. Nanopoulos, Nucl. Phys. B135 (1978) 66;
S. Weinberg, Phys. Lett. 82B (1979) 387;
M. Veltman, Acta Physica Polonica, B12 (1981) 437.
- 4) E. Witten, Nucl. Phys. B188 (1981) 513;
S. Dimopoulos and H. Georgi, Nucl. Phys. B193 (1981) 150;
N. Sakai, Zeit. für Physik, C11 (1981) 153;
R.K. Kaul, Phys. Lett. 109B (1982) 19.
- 5) A. Masiero, D.V. Nanopoulos, K. Tamvakis and T. Yanagida, CERN preprint TH-3298 (1982), to appear in Phys. Lett. B.
- 6) For a review, see:
D.V. Nanopoulos, in Ecole d'Eté de Physique des Particules, Gif-sur-Yvette (1980) (IN2P3, Paris 1980), p. 1.
- 7) D.V. Nanopoulos and K. Tamvakis, Phys. Lett. 110B (1982) 449.
- 8) M. Srednicki, Princeton preprints (1982).
- 9) D.V. Nanopoulos, K. Olive and K. Tamvakis, CERN preprint TH-3278 (1982), to appear in Phys. Lett. B.
- 10) D.V. Nanopoulos and K. Tamvakis, CERN preprint TH-3255 (1982), to appear in Phys. Lett. B.
- 11) A. Masiero, D.V. Nanopoulos, K. Tamvakis and T. Yanagida, MPI preprint MPI-PAE/PTh 29/82 (1982), to appear in Phys. Lett. B;
Y. Igarashi, T. Kubo and S. Sakakibara, Dortmund preprint DO-TH-82/09 (1982).
- 12) D.V. Nanopoulos and K. Tamvakis, Phys. Lett. 113B (1982) 151.
- 13) L. Girardello, M. Grisaru and P. Salomonson, Nucl. Phys. B178 (1981) 331.
- 14) S.B. Treiman and F. Wilczek, Phys. Lett. 95B (1980) 222.
- 15) R. Barbieri, D.V. Nanopoulos and D. Wyler, Phys. Lett. 103B (1981) 433; *ibid.* 106B (1981) 303.
- 16) J. Ellis, M.K. Gaillard and D.V. Nanopoulos, Phys. Lett. 80B (1979) 360.
- 17) J. Ellis, D.V. Nanopoulos and S. Rudaz, CERN preprint TH-3199 (1981).
- 18) H. Georgi and D.V. Nanopoulos, Phys. Lett. 82B (1979) 95.

- 19) S. Weinberg, Harvard preprint HUTP-81/A04 (1981).
- 20) N. Sakai and T. Yanagida, Nucl. Phys. B197 (1982) 533.
- 21) J. Ellis and D.V. Nanopoulos, Phys. Lett. 110B (1982) 44.
- 22) R. Barbieri, D.V. Nanopoulos and A. Masiero, Phys. Lett. 98B (1981) 191;
A. Masiero and R.N. Mohapatra, Phys. Lett. 103B (1981) 343.
- 23) A. Masiero and T. Yanagida, Phys. Lett. 112B (1982) 336.
- 24) A. Masiero, J.F. Nieves and T. Yanagida, MPI preprint MPI-PAE-PTh 26/82.

FIGURE CAPTIONS

- Figure 1 : Two-loop cut graph contributing to the CBA.
- Figure 2 : Diagrams giving rise to operators of dimension five.
- Figure 3 : Diagrams giving rise to operators of dimension six.
- Figures 4&5: Two-loop cut graphs contributing to the CBA in the model given in Eq. (15).
- Figure 6 : Three-body decay for superheavy fermions which cannot decay directly into light particles through two-body decays. See Ref. 22).
- Figure 7 : Two-body decays for superheavy fermions which couple directly to the light matter. Examples are given in Ref. 23).
- Figure 8 : Φ mediated proton decay.
- Figure 9 : B+L-conserving decay of a nucleon with a lepton and a kaon in the final state.

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