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Heavy Majorana leptons and cosmological baryon excess

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Cosmological baryon excess is used to probe how the left-right symmetry should be broken in a basically symmetric grand unified theory. In the SO(10) model, the contribution from an intermediate state containing the chiral partner of a light ordinary neutrino cancels much of the usual baryon excess produced by a heavy-Higgs-boson decay unless the partner's mass is $\gtrsim 10^{10}$ GeV. A mechanism of soft CP violation is also mentioned.

I. INTRODUCTION

Grand unified theories (GUT's) that unite the electroweak and strong interactions of fundamental particles have been a subject of extensive studies. In particular, the simplest and the most economical model of GUT's based on the SU(5) group has been discussed in greatest detail.[1] There are, however, several reasons, either based on sound theoretical prejudices or on practical considerations, to suspect that this simple SU(5) model is not the whole story. Two major theoretical arguments against the SU(5) model are its left-right (LR) asymmetric nature and the exact conservation of $B-L$, namely, baryon minus lepton number. We regard the conservation of $B-L$ as accidental and quite arbitrary because this global conservation law is not tied up with the fundamental gauge principle. In fact, $B-L$ can be violated[2] even in the SU(5) model by arbitrarily introducing a new Higgs multiplet, $\underline{15}$ or $\underline{10}$ of SU(5). We reject this solution because it is highly artificial. A more practical problem present in the minimal SU(5) model is that the cosmological baryon excess computed in this case is found to be several orders of magnitude too small[3] in comparison to cosmological observation. Again, this difficulty is overcome[3] by introducing more $\underline{5}$'s of Higgs bosons, but this solution is not attractive either.

In view of the problems just mentioned, we believe it worthwhile to construct and work out physical consequences of a simple and appealing model that overcomes these theoretical objections against the SU(5) model. Our model is based on the SO(10) group[4] that includes the SU(5) as a subgroup. One advantage of SO(10) grand unified theories is that by imposing a discrete reflection symmetry, the model is made C, P, and CP conserving prior to a spontaneous symmetry breakdown. All these discrete symmetries in this case are broken only softly, namely, spontaneously along with the gauge symmetry. This soft violation has a number of interesting consequences. In general, the gauge hierarchy of grand unification implies a hierarchy of soft CP violation. Namely, a number of relative phase factors of Higgs-field vacuum expectation values that characterize CP violation may arise from different Higgs multiplets underlying different energy scales of the gauge hierarchy. This seems realistic in the sense that we have at least two different kinds of effects of CP violation observed in nature; the cosmological baryon asymmetry and the CP violation in K decay. On the other hand, this hierarchy is a great obstacle against an attempt to explain the sign of the baryon asymmetry relative to that in the K decay. Perhaps the most interesting result in a minimal version of this type of model is on neutrino masses and mixing, which was already described in a short communication.[5] By minimizing the number of Higgs multiplets in the SO(10) model, we have obtained a testable prediction on neutrino oscillation: We have found that only one neutrino mass dominates and that the mixing matrix of leptonic weak eigenstates coincides with that of quarks. On this neutrino oscillation we refer to our previous paper;[5] and we shall not discuss it much in this paper.

The problem of the cosmological baryon excess in SO(10) models has a new aspect that was not encountered in the SU(5) model. From the very LR-symmetric nature of SO(10) models, neutral leptons are described by four-component objects unlike two-component neutrinos in the SU(5) model. An important question then arises whether these are lumped together as massive Dirac spinors or split into two types of Majorana leptons, each with two components. Note that Majorana mass terms explicitly violate $B-L$. It is found from arguments of cosmological baryon production that the existence of new heavy (of masses $\gtrsim 10^{10}$ GeV) Majorana leptons is definitely favored over light Dirac neutrinos. Heavy Majorana leptons as massive as ours imply an apparent lack of the left-right symmetry prior to the grand unification symmetry. Existence of the heavy Majorana leptons in turn enhances expectation of neutrino oscillation, as recently discussed.[6,7]

This paper is organized as follows. In Sec. II we describe the main ingredients of our SO(10) model. In Sec. III we compute the cosmological baryon excess produced by decays of various heavy particles: gauge bosons, Higgs bosons, and heavy Majorana particles. We stress special features that appear in the SO(10) model, in particular, a crucial role played by the heavy Majorana lepton. The final section is devoted to a summary and discussions.

II. SO(10) MODEL

A basic set of fermions used in this model forms a family of irreducible 16-plet of the SO(10) group ($i=1,\ldots,$ number of families),

$$\psi_{iL}=(N^c,e^c,u,d,u^c,d^c,e,\nu)_{iL}\,. \tag{1}$$

Here, u_i and d_i, with three colors omitted for simplicity, denote quarks of charges $\frac{2}{3}$ and $-\frac{1}{3}$, respectively, and e_i is a charged lepton. The superscript c means the usual charge conjugation and $L=(1-\gamma_5)/2$. We deliberately separated a four-component neutral lepton into two chiral components, ν_L and $(N^c)_L$, or, equivalently, ν_L and N_R, anticipating the result that they eventually correspond to light ($\lesssim$ eV) and heavy ($\gtrsim 10^{10}$ GeV) Majorana particles[6,7] after a spontaneous symmetry breakdown. Under the SU(5) subgroup the $\underline{16}$ is decomposed into

$$\begin{aligned}\psi_{iL}(\underline{16})&=\xi_{iL}[\underline{10}]+(\eta_{iR})^c[\underline{5}^*]+(N_{iR})^c[\underline{1}]\\&=\begin{bmatrix}e^c & u & u^c\\ & d & \end{bmatrix}_{iL}+\begin{bmatrix}\nu & d^c\\ e & \end{bmatrix}_{iL}+(N^c)_{iL}\,.\end{aligned} \tag{2}$$

Hereafter we use parentheses and square brackets to indicate dimensionalities of the SO(10) and SU(5) groups, respectively. ξ and η in Eq. (2) are the usual fermions that appear already in the Georgi-Glashow SU(5) model,[1] and a new SU(5) singlet N was introduced to complete the SO(10) representation. Another useful subgroup of SO(10) is SO(6) × SO(4), which is locally equivalent to SU(4) × SU(2) × SU(2)′. In terms of this subgroup,

$$\begin{aligned}\psi_{iL}(\underline{16})&=\{\underline{4},\underline{2},\underline{1}\}_{iL}+\{\underline{4}^*,\underline{1},\underline{2}^*\}_{iL}\\&=\begin{pmatrix}u_1 & u_2 & u_3 & \nu\\ d_1 & d_2 & d_3 & e\end{pmatrix}_{iL}+\begin{pmatrix}u_1^c & u_2^c & u_3^c & N^c\\ d_1^c & d_2^c & d_3^c & e^c\end{pmatrix}_{iL}\,.\end{aligned} \tag{3}$$

The subgroup SU(4) can be identified with the Pati-Salam[8] unification group if the fractional charge assignment of quarks is adopted instead of the integral one in the Pati-Salam model, and also the two subgroups of SU(2) can be interpreted as the left-right-symmetric $SU(2)_L \times SU(2)_R$ group.[9] This $SU(4) \times SU(2)_L \times SU(2)_R$ may not be a good intermediate gauge symmetry, but this subgroup is still useful as a classification scheme.

A grand unified model is never complete unless a Higgs-boson mechanism is explicitly specified. Two possible symmetry-breaking patterns are as follows:

$$\begin{tikzcd}[column sep=small] & SU(5) \arrow[dr, "\langle\underline{45}\rangle"] & & \\ SO(10) \arrow[ur, "\langle\underline{16}\rangle"] \arrow[dr, "\langle\underline{210}\rangle \text{ or } \langle\underline{54}\rangle"'] & & SU(3)\times SU(2)\times U(1) \arrow[r, "\langle\underline{10}\rangle"'] & SU(3)\times U(1)\,. \\ & SU(4)\times SU(2)_L\times SU(2)_R \arrow[ur, "\langle\underline{45}\rangle"'] & & \end{tikzcd}$$

(4a)

(4b)

We have shown vacuum expectation values of Higgs multiplets relevant to respective symmetry breakdowns in the simple case. We shall have more to say about the pattern of symmetry breakdown when we discuss the cosmological baryon excess in Sec. III. With irreducible families of Eq. (1), mass terms and Yukawa couplings to Higgs bosons are deduced from the bilinear forms of

$$(\overline{\psi}_{iL})^c\psi_{jL}=\psi_{iL}^T C\psi_{jL}=\psi_{jL}^T C\psi_{iL}\,. \tag{5}$$

The last equality was derived from the Fermi statistics (anticommutation of ψ's). As to the SO(10) content, irreducible representations contained here are

$$\underline{16}\times\underline{16}=\underline{10}+\underline{120}+\underline{126}\,. \tag{6}$$

We shall only write the Yukawa coupling of the simplest Higgs multiplet $\underline{10}$, $\varphi_a(\underline{10})$ ($a=1,\ldots,$ number of $\underline{10}$'s) as follows:

$$L_Y=(\overline{\psi}_{iL})^c f_{ij}^a\varphi_a(\underline{10})\psi_{jL}+\text{H.c.}, \tag{7}$$

where the Yukawa coupling matrix f^a is symmetric by Eq. (5) and by the symmetry of the Clebsch-Gordan coefficients of $\underline{10}$. In order to see a relation to the SU(5) model, it is convenient to introduce complex fields

$$\chi_{ak}\equiv(\varphi_{a2k}+i\varphi_{a2k-1})/\sqrt{2}\quad(k=1\text{–}5)\,, \tag{8}$$

which transform as $\underline{5}$ under the SU(5) subgroup. In terms of this $[\underline{5}]$, and $\xi_L[\underline{10}]$, $\eta_R[\underline{5}]$, and $N_R[\underline{1}]$ of Eq. (2), Eq. (7) can be rewritten

$$\begin{aligned}L_Y/\sqrt{2}=[&\overline{\xi}_{iL}\eta_{jR}f_{ij}^{a*}-(\overline{\xi}_{iL})^c\xi_{jL}f_{ij}^a\\&+\overline{\eta}_{iR}(N_{jR})^c f_{ij}^a]\chi_a+\text{H.c.}.\end{aligned} \tag{9}$$

Definition of f^a in Eqs. (7) and (9) is not precise, but $f^{(a)}$ in the notation of Eq. (2.9) of Ref. 10 is equal to $\sqrt{2}f^{a*}$ here and $h^{(a)}$ in Eq. (2.9) is equal to $-\sqrt{2}f^a$ here. At least two $\underline{10}$'s of Higgs are found necessary[5] to get a reasonable mass spectrum of quarks and leptons.

An interesting mechanism of producing large Majorana masses of $\gtrsim 10^{10}$ GeV for the chiral component N was proposed by Witten[7] and analyzed by us[5] in more detail in connection with neutrino masses and neutrino mixing. The result, with respect to ordinary neutrinos, amounts to introducing a neutrino mass matrix of Majorana type

$$M_\nu = M_u M_N^{-1} M_u\,, \tag{10}$$

where M_N is a large mass matrix of N and M_u that of quarks of charge $\frac{2}{3}$. For the implications of Eq. (10) on the phenomena of neutrino oscillation, we refer to our previous paper.[5] This mechanism uses the SO(10) representations of $\underline{16}$, $\underline{45}$, and $\underline{10}$ as Higgs multiplets.

Finally, we point out a possibility of implementing a soft CP violation. By demanding a discrete reflection symmetry R, in addition to symmetry operations of the continuous Lie group SO(10), the model can be made C, P, and CP conserving prior to a symmetry breakdown. For this purpose the spinorial $\underline{16}$ must either be doubled by adding a mirror $\overline{\underline{16}}$ or be made self-conjugate, both in a form of reducible $\underline{32}$. The simplest model of this kind contains Higgs multiplets of two $\underline{16}$'s, one $\underline{45}$, and two $\underline{10}$'s and three families of self-conjugate fermions of $\underline{32}$. A relative phase of two vacuum expectation values (VEV's) of $\underline{16}$ gives CP violation even in the SU(5) symmetry limit, which then becomes a major source of CP violation in the cosmological baryon excess. The observed CP violation, as measured in the K decay, is induced by having a relative phase of two VEV's of 10 such that $\arg(\langle\chi_1\rangle/\langle\chi_2\rangle) \equiv \delta \neq 0$ or $\pi/2$ in the notation of Eq. (8). Among the Yukawa coupling matrices of Eq. (7), f^1 is real and f^2 is purely imaginary by requiring the soft CP violation. The Kobayashi-Maskawa phase in the case of three generations is related to the relative phase δ, but this model of soft CP violation can in principle be distinguished from the quark mixing model if the number of generations exceeds three.

III. COSMOLOGICAL BARYON EXCESS

This subject in general and specific computations in SU(5) models have been discussed in many recent papers.[11,3] A key quantity in this problem is the net baryon number ΔB produced when a pair of heavy particles and its antiparticles (if they are distinguishable) decay via baryon- and CP-nonconserving interactions. Possible candidates of such heavy particles in SO(10) models are gauge bosons (X and their relatives denoted in general by X_v), colored Higgs bosons (X_s), and heavy Majorana leptons (N). We shall compute the asymmetry ΔB produced in decay processes of these heavy particles. This asymmetry is related to the observed cosmological ratio of baryon to photon number N_B/N_γ by a calculable factor[11] of ~ 100; thus ΔB must be $\sim 10^{-7\pm1}$ in order to explain N_B/N_γ. For simplicity, we shall exhibit results in the SO(10) model that contains the Yukawa coupling of $\underline{10}$ alone. We shall not discuss another important aspect[12] of kinematical constraint in this paper.

A. Gauge-boson decay

Gauge bosons that can decay to a number of channels of difermions with different baryon numbers have the following $SU(3)_c \times SU(2)_{WS}$ (where WS denotes Weinberg-Salam) quantum numbers,[13]

$$X_v(-\tfrac{1}{3}, -\tfrac{4}{3})\,, \quad (\underline{3}, \underline{2})\,, \tag{11}$$

$$X_v'(\tfrac{2}{3}, -\tfrac{1}{3})\,, \quad (\underline{3}, \underline{2})\,. \tag{12}$$

Their antiparticles have conjugate quantum numbers. Fractional numbers in the first parentheses are electric charges of these bosons. Both types of X bosons are present in SO(10) models, while only X_v is present in SU(5) models. Gauge bosons are known to be rather ineffective[3] in producing the baryon asymmetry ΔB, and the only known mechanism of Ref. 10 uses Higgs-boson intermediate states. If only X_v or X_v' contributes to ΔB in the SO(10) model that contains the Yukawa coupling of $\underline{10}$ alone, then the result is basically the same[10] as in SU(5) models,

$$\Delta B \simeq O\left(\frac{1}{12\pi}\,\mathrm{tr} f f^\dagger\right)\,, \tag{13}$$

for the three families of fermions. We have ignored mass differences of $O(1)$ and assumed a maximal CP violation. A nice feature of the SO(10) model is that there is a good reason, as already mentioned, for introducing an extra $\underline{10}$ of Higgs boson owing to the fermion mass spectrum. Existence of the second $\underline{10}$ is necessary to obtain the asymmetry of order of Eq. (13), as seen from detailed discussions in Ref. 10. Moreover, the corresponding, second Weinberg-Salam type of Higgs doublet must be superheavy[10] to give a sizable ΔB, which is presumably nice because this completely suppresses[14] effects of simulated flavor-changing neutral currents due to the second Higgs doublet. Unfortunately, if CP violation is soft or LR symmetry is imposed on the Yukawa coupling, then this mechanism of baryon production does not work because the coupling matrix f is real and a cancellation takes place for ΔB.

B. Higgs-boson decay

A dominant class of diagrams that contribute to the baryon asymmetry ΔB is depicted in Fig. 1, where appropriate final-state and discontinuity cuts should be made. Again, existence of many $\underline{10}$'s is a blessing[15] in this case, too. Relevant heavy Higgs bosons contributing to ΔB have the following, definite $SU(3)_c \times SU(2)_{WS}$ quantum numbers

$$X_s(-\tfrac{1}{3}), \quad (\underline{3},\underline{1}), \tag{14}$$

in the SO(10) model containing the Yukawa coupling of $\underline{10}$ alone. The asymmetry induced by the decays $(X_s)_a \to \overline{q}\overline{q}$ and ql is given by

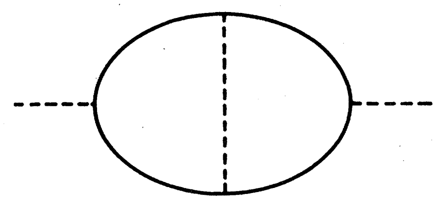

FIG. 1. Higgs-boson decay $X_s \to \overline{q}\overline{q}$ and ql contributing to the baryon excess. Dotted and solid lines represent heavy Higgs bosons and fermions, respectively.

$$(\Delta B)_a = -\frac{1}{8\pi}\left[\frac{7}{8}\operatorname{tr}(f_a^\dagger f_a) + \sum_i \frac{1}{8}\left(1-\frac{\mu_i^2}{m_a^2}\right)^2 \operatorname{tr}(f_a^\dagger f_a \rho_i \rho_i^\dagger)\right]^{-1}$$
$$\times \operatorname{Im}\sum_{bi}\operatorname{tr}\left\{f_b^\dagger f_a f_b^\dagger f_a\left[\rho_i\rho_i^\dagger g\left(\frac{m_b^2}{m_a^2},\frac{\mu_i^2}{m_a^2}\right) - g\left(\frac{m_b^2}{m_a^2},0\right)\right]\right\}, \tag{15}$$

where f_a is the same Yukawa coupling matrix as in Eqs. (7) and (9), denoted there by the matrix f^a. ρ_i is a projection matrix onto a Majorana mass eigenstate N_i, which is symmetric and in general complex, but not Hermitian. By a straightforward calculation it is shown that

$$g(x,y) = \left[1 - y - x\ln\left(\frac{1+x-y}{x}\right)\right]\theta(1-y). \tag{16}$$

The step function $\theta(1-y)$ here ensures existence of a discontinuity for $(X_s)_a \to \overline{u}\overline{d}$ via intermediate states ql and dN. As clearly seen from the explicit computation of Eq. (15), there is a partial cancellation of contributions that come from intermediate states involving N and ν. Indeed, if all masses of Majorana leptons μ_i's $\ll m_a$ of the parent Higgs boson, then this cancellation is almost complete and ΔB nearly vanishes. A basic reason for this cancellation is that under the parity operation the Yukawa coupling f gets transformed to f^*.

In order to estimate Eq. (15), let us assume that all Higgs-boson masses are the same, $m_a = m$. Then, up to the Majorana mass $\lesssim 0.6m$, numerically,

$$(\Delta B)_a \simeq \frac{1}{16\pi}[\operatorname{tr}(f_a f_a^\dagger)]^{-1}\operatorname{Im}\sum_{bi}\operatorname{tr}[(f_b^\dagger f_a f_b^\dagger f_a \rho_i \rho_i^\dagger \mu_i^2)/m^2]. \tag{17}$$

With a maximal CP violation and three families of fermions, a crude estimate of the right-hand side (RHS) is given by

$$(\mathrm{RHS}) \sim O(\tfrac{1}{16}\,\alpha_{GU}\,\mu^2 m_t^2/(9m^2 m_W^2)). \tag{18}$$

A factor of 9 was put in to take into account the renormalization effect of the top-quark mass.[16] With a reasonable guess of the gauge coupling $\alpha_{GU} \sim \frac{1}{50}$, we have a range of mass ratio

$$\mu/m \sim 9\times 10^{-2} \times 25\ \mathrm{GeV}/m_t \tag{19}$$

with uncertain factors of (3–0.3) derived in order to get $\Delta B \sim 10^{-7\pm 1}$. A mass m of colored Higgs boson must be greater than $\sim 10^{12}$ GeV in order to suppress the nucleon decay. This order of magnitude given by Eq. (19) then falls within the interesting range[7] of a Majorana-lepton mass to make the related neutrino mass ~ 10–10^{-4} eV. The necessity of a large Majorana-lepton mass substantiates the general argument[17] claiming that restoration of the left-right symmetry prior to the grand unification symmetry is excluded by the baryon excess ΔB. In our case, the left-right symmetry prior to grand unification would mean other things almost massless neutral leptons, in contradiction to Eq. (19). A mass at least as heavy as $\sim 10^{10}$ GeV seems necessary for the Majorana lepton N.

Finally, in the case of soft CP violation, the Yukawa couplings f that appear in Eq. (15) can effectively be treated as real. Thus, at least one of the projection matrices ρ_i onto Majorana-lepton mass eigenstates must contain an imaginary part. This is possible only if there are more than two $\underline{16}$'s of Higgs multiplets and, furthermore, if nonvanishing relative phases of SU(5)-singlet VEV's of $\underline{16}$ are demanded by minimization of the Higgs potential.

C. Heavy-Majorana-lepton decay

Two possible mechanisms of baryon production via N decay in a more general class of GUT's have been suggested in our recent publication.[18] In the case of the SO(10) model of this paper, we obtain more specific results, which we shall report here. For the three body decays $N_i \to qqq$ and $N_i \to \bar{q}\bar{q}\bar{q}$, the produced baryon asymmetry is given by

$$(\Delta B)_i = -(4\pi^4)^{-1}\left[\sum_a \mathrm{tr}(\rho_i\rho_i^\dagger f_a^\dagger f_a)\right]^{-1} 3\,\mathrm{Im}\sum_{abcj}\mathrm{tr}(\rho_i f_a f_b^\dagger f_c f_b^\dagger \rho_j^\dagger f_a^\dagger f_c) h\left(\frac{m^2}{\mu_i^2}, \frac{\mu_j^2}{\mu_i^2}\right)\theta(\mu_i^2 - \mu_j^2), \tag{20}$$

where a common mass m was assumed for several Higgs bosons of X_s which appear in relevant diagrams of Fig. 2. The function $h(x,y)$ can be computed only numerically from the discontinuity of Fig. 2 and typical numerical results are shown in Table I for representative values of the mass ratios m^2/μ_i^2 and μ_j^2/μ_i^2 of Eq. (20). This mechanism is dominant only for a relatively light Majorana lepton, namely, for the case of a Majorana-lepton mass $\mu <$ Higgs-boson mass m. The maximal absolute value of $h(x,y)$ is around 3×10^{-3} as seen from Table I. A crude estimate then gives, with a maximal CP violation and three generations of fermions,

$$|(\Delta B)_i| \lesssim 3\times 10^{-6}\alpha_{\mathrm{GU}}^2\left(\frac{m_t}{m_W}\right)^4. \tag{21}$$

This appears too small to give a reasonable value of the baryon-to-photon ratio N_B/N_γ.

Another possibility[18] of baryon production manifests itself when a Majorana lepton N_i is heavier than some of the X_s bosons since the decay $N_i \to X_s\bar{q}$ or $\bar{X}_s q$ followed by X_s decay becomes possible. A diagram of Fig. 3 gives a contribution to ΔB given by

$$(\Delta B)_i = -\frac{12}{35}(8\pi)^{-1}\sum_{abj}[\mathrm{tr}(\rho_i\rho_i^\dagger f_a^\dagger f_a)]^{-1}\,\mathrm{Im}\,\mathrm{tr}(\rho_i^\dagger f_b^\dagger f_a \rho_j f_b f_a^\dagger)$$
$$\times\left(1-\frac{m^2}{\mu_i^2}\right)^2\left[\frac{3}{5}\left(1-\frac{m^2}{\mu_i^2}\right)^2+\frac{2}{5}\right]^{-1} k\left(\frac{m^2}{\mu_i^2}, \frac{\mu_j^2}{\mu_i^2}\right), \tag{22}$$

where $-\frac{12}{35}$ is the fractional baryon number produced by successive decays of N and X_s if all mass differences are ignored. For simplicity, we assumed a common Higgs-boson mass m for various X_s. From a straightforward computation,

$$k(x,y) = \sqrt{y}\left[1+\frac{1+y-2x}{(1-x)^2}\ln\left(\frac{y-x^2}{1+y-2x}\right)\right]\theta(1-x). \tag{23}$$

For a simple estimate let us take a common mass μ for all Majorana leptons N. Assuming a maximal CP violation, one obtains a crude estimate with three generations of fermions,

$$\Delta B \simeq 5\times 10^{-3}\alpha_{\mathrm{GU}}\left(\frac{m_t}{m_W}\right)^2 Y\left(\frac{m^2}{\mu^2}\right), \tag{24}$$

where

$$Y(x) = (1-x)^2[\tfrac{3}{5}(1-x)^2+\tfrac{2}{5}]^{-1}$$
$$\times\left[1+\frac{2}{1-x}\ln\left(\frac{1+x}{2}\right)\right]\theta(1-x). \tag{25}$$

Numerically, $|Y(x)|$ is monotonically decreasing and

$$0.4 > |Y(x)| > 6\times 10^{-4} \quad \text{when } 0<x<0.9.$$

Thus, with $m_t = 25$ GeV, $\Delta B = 10^{-7}$ corresponds to the case when $m \sim 0.9\mu$.

IV. SUMMARY AND DISCUSSIONS

From computations of the baryon asymmetry in the previous section we have deduced one important conclusion: The left-right symmetry must be broken above, or nearly at, the grand unification scale of $\sim 10^{15}$ GeV in order to obtain a reasonable size of the baryon asymmetry ΔB. In the case of

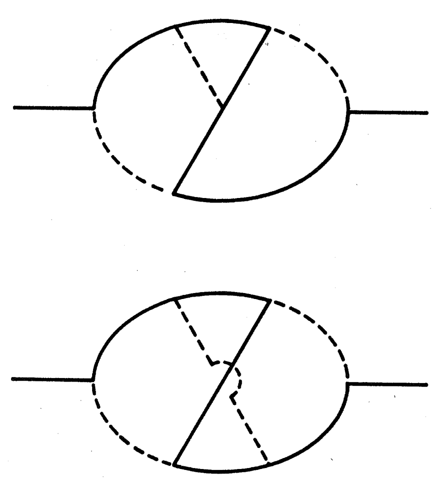

FIG. 2. Heavy-Majorana-lepton decay $N \to qqq$ and and $\bar{q}\bar{q}\bar{q}$ contributing to the ΔB. Dotted and solid lines represent heavy Higgs bosons and fermions, respectively.

TABLE I. Numerical values of the function h of Eq. (20). m, μ, and μ' here denote masses of exchanged Higgs, parent Majorana, and intermediate Majorana particles, respectively.

$m \backslash \mu'$	$0.1\ \mu$	$0.3\ \mu$	$0.5\ \mu$
μ	1×10^{-3}	3×10^{-3}	2×10^{-3}
2μ	1×10^{-5}	2×10^{-5}	8×10^{-6}
5μ	4×10^{-8}	6×10^{-8}	3×10^{-8}

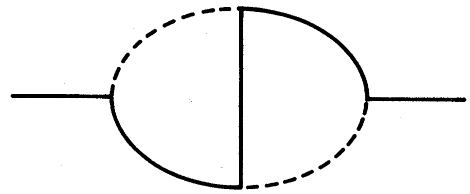

FIG. 3. Heavy-Majorana-lepton decay $N\to X_s\,\bar{q}$ and $\bar{X}_s q$ contributing to ΔB when followed by X_s $(\bar{X}_s)\to\bar{q}\,\bar{q}$ (qq). Dotted and solid lines represent heavy Higgs bosons and fermions, respectively.

the gauge-boson decay the Yukawa coupling of Higgs 10 becomes real by imposing LR symmetry and the gauge-boson decay is not effective for generation of the asymmetry ΔB. Thus, LR symmetry must be broken explicitly, or other Higgs muliplets 120 or 126 must be important to give a sizable ΔB in this case. In the other cases of Higgs-boson and neutral-lepton decays, violation of the LR symmetry is realized by a hierarchy of masses of the ordinary neutrino and its chiral partner N, which is effective to give a reasonable ΔB. Both cases support the hierarchical pattern of symmetry breakdowns as shown in Eq. (4a). We believe that this conclusion for the lack of the left-right symmetry is much more general, and in fact its validity can be checked in more general SO(10) models in which Yukawa couplings of 120 and 126 in the sequence of Eq. (6) are also present. We also mention that a general argument against the left-right symmetry has been advanced.[17] Our specific computations reveal much more than this general argument indicates. We can learn quantitatively how badly the left-right symmetry must be broken spontaneously in a basically symmetric theory such as the SO(10) model. For example, Eq. (19) or Eq. (24) tells how large a mass of the heavy Majorana lepton must be if the Higgs-boson or the heavy-Majorana-lepton decay is a dominant cause of baryon production. This in turn makes it possible to estimate a size of neutrino masses through a formula such as Eq. (10). Thus, consideration of the cosmological baryon excess gives important constraints on the theory of grand unification.

We have also mentioned a way to implement a theory of soft CP violation in the context of SO(10) models. There are many to be studied in this type of soft model. An interesting problem related to the cosmological baryon is whether this type of soft model of CP violation leads to a domain structure[19] of matter and antimatter on a cosmological scale. We hope to return to these problems in a future communication.

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[1] H. Georgi and S. L. Glashow, Phys. Rev. Lett. 32, 438 (1974); A. J. Buras, J. Ellis, M. K. Gaillard, and D. V. Nanopoulos, Nucl. Phys. B135, 66 (1978).
[2] F. Wilczek and A. Zee, Phys. Lett. 88B, 311 (1979).
[3] S. M. Barr, G. Segre, and H. A. Weldon, Phys. Rev. D 20, 2494 (1979); T. Yanagida and M. Yoshimura, Nucl. Phys. B168, 534 (1980); A. Yildiz and P. Cox, Phys. Rev. D 21, 906 (1980); D. V. Nanopoulos and S. Weinberg, *ibid.* 20, 2484 (1979).
[4] H. Georgi, in *Particles and Fields—1974*, proceedings of the 1974 Meeting of the APS Division of Particles and Fields, Williamsburg, edited by C. E. Carlson (AIP, New York, 1975); H. Fritzsch and P. Minkowski, Ann. Phys. (N.Y.) 93, 193 (1975); M. S. Chanowitz, J. Ellis, and M. K. Gaillard, Nucl. Phys. B128, 506 (1977); H. Georgi and D. V. Nanopoulos, *ibid.* B155, 52 (1979).
[5] T. Yanagida and M. Yoshimura, Phys. Lett. 97B, 99 (1980).
[6] M. Gell-Mann, P. Ramond, and R. Slansky, in *Supergravity*, edited by D. Z. Freedman and P. van Nieuwenhuizen (North-Holland, Amsterdam, 1979).
[7] E. Witten, Phys. Lett. 91B, 81 (1980).
[8] J. C. Pati and A. Salam, Phys. Rev. D 8, 1240 (1973).
[9] J. C. Pati and A. Salam, Phys. Rev. D 10, 275 (1974); R. N. Mohapatra and J. C. Pati, *ibid.* 11, 566 (1975); 11, 2558 (1975); G. Senjanovic and R. N. Mohapatra, *ibid.* 12, 1502 (1975).
[10] T. Yanagida and M. Yoshimura, Ref. 3.
[11] M. Yoshimura, Phys. Rev. Lett. 41, 281 (1978); 42, 746(E) (1979); S. Dimopoulos and L. Susskind, Phys. Rev. D 18, 4500 (1978); D. Toussaint, S. B. Treiman, F. Wilczek, and A. Zee, *ibid.* 19, 1036 (1979); S. Weinberg, Phys. Rev. Lett. 42, 850 (1979).

[12]M. Yoshimura, Phys. Lett. 88B, 294 (1979); M. Honda and M. Yoshimura, Prog. Theor. Phys. 62, 1704 (1979); E. W. Kolb and S. Wolfram, Nucl. Phys. B172, 224 (1980); J. N. Fry, K. A. Olive, and M. S. Turner, Phys. Rev. D 22, 2953; 22, 2977 (1980).
[13]S. Weinberg, Phys. Rev. Lett. 43, 1566 (1979); F. Wilczek and A. Zee, *ibid.* 43, 1571 (1979).
[14]H. Georgi and D. V. Nanopoulos, Phys. Lett. 82B, 95 (1979).
[15]Yildiz and Cox, Ref. 3; Nanopoulos and Weinberg, Ref. 3.
[16]A. J. Buras *et al.*, Ref. 1.
[17]V. A. Kuzmin and M. E. Shaposhnikov, Phys. Lett. 92B, 115 (1980).
[18]T. Yanagida and M. Yoshimura, Phys. Rev. Lett. 45, 71 (1980).
[19]G. Senjanovic and F. W. Stecker, Phys. Lett. 96B, 285 (1980) and references therein; I. Yu. Kobsarev, L. B. Okun, and Ya. B. Zeldovich, Phys. Lett. 50B, 340 (1974).