sponsible for the diffusion process as suggested by the similarity of the activation energies with that found for ^{3}H in ice.[16]

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Cosmological and Astrophysical Implications of Heavy Majorana Particles

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It is pointed out that heavy Majorana particles contemplated in certain grand unified theories may explain the cosmological baryon excess. Induced neutrino mixing can give oscillation lengths consistent with reactor and accelerator experiments, but capable of explaining the solar-neutrino "puzzle."

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It is somewhat puzzling that neutrinos appear massless in sharp contrast to all other quarks and charged leptons. This may simply be a reflection of extremely large masses of their chiral partners (Majorana type), in which case a sizable mixing comparable to usual mass terms is allowed between neutrinos and heavy Majorana particles. This situation can happen[1] in a number of grand unified models, including some SO(10) models,[2] that are more left-right symmetric and therefore aesthetically more appealing than the minimal SU(5) model.[3] In this note we shall demonstrate two possible implications of these heavy Majorana particles: generation of the cosmological baryon asymmetry and a large oscillation length of neutrinos that may explain the solar-neutrino "puzzle."[4] In the picture that we shall develop below, these two separate issues in cosmology and astrophysics are correlated and have a common origin in extremely large masses ($\gtrsim 10^{10}$ GeV) of Majorana particles.

In order to clarify our main points we shall take as an illustration a simple model that extends the minimal SU(5) model[3] by simply adding SU(5)-singlet Majorana particles N_i ($i = 1 \sim$ number of generations). This particular example should serve as a prototype of more complicated situations like SO(10) models.[2] The two-component fields N_i do not couple to SU(5) gauge bosons, but may have Yukawa couplings to the five-dimensional Higgs boson $H(\underline{5})$ usually introduced in SU(5) models. Mass terms and Yukawa couplings to $H(\underline{5})$ of fundamental fermions, $\varphi_i(\underline{5})$, $\psi_i(\underline{10})$, and $N_i(\underline{1})$, are given by

$$L_{m,\mathrm{Y}} = \bar{\tilde{N}}(\underline{1}) M N_R(\underline{1}) + [\bar{\psi}(\underline{10}) f \varphi_R(\underline{5}) + \bar{\tilde{\psi}}(\underline{10}) h \psi_L(\underline{10}) + \bar{\varphi}(\underline{5}) \lambda^\dagger \tilde{N}_R(\underline{1})] H(\underline{5}) + (\mathrm{H.c.}). \tag{1}$$

Here, M, f, h, and λ are matrices of generation indices and $L, R = \frac{1}{2}(1 \mp \gamma_5)$ are usual chiral projections. The tilde in Eq. (1) means charge conjugation. This form (1) has the most general structure allowed by the SU(5) gauge invariance in the minimal Higgs system that contains one $H(\underline{5})$ and one adjoint $(\underline{24})$ representation.

By use of appropriate unitary transformations, the matrices M and f in Eq. (1) can be made real and diagonal, but h and λ are in general complex except that $h^T = h$ (symmetric) because of the two $\psi(\underline{10})$'s in (1). These complex parameters are the only sources of CP-invariance violation in this model. After spontaneous symmetry breaking of the SU(5) a neutral component of $H(\underline{5})$ acquires a nonvanishing vacuum expectation value v so that all charged fermions become massive and mass terms of neutral leptons are given by

$$\bar{\tilde{N}} M N_R + v \bar{\nu} \lambda^T N_R + (\mathrm{H.c.}). \tag{2}$$

Off-diagonal terms $v\lambda_{ij}$ in this equation give rise to a mass mixing between Majorana (N) and two-component Weyl (ν) fields. Since we have in mind a grandunification with a group larger than SU(5), it is not unreasonable to assume M to be of order $\approx 10^{15}$ GeV and $v\lambda$ to be of order $\lesssim 10^2$ GeV, typical of ordinary quarks and leptons. After diagonalization of Eq. (2) two types of Majorana particles emerge as mass eigenstates, one denoted by N_i' of masses $\approx M$ and the other denoted by ν_i' of masses $\approx v^2\lambda^2/M$. Neutrino (ν_i') masses are, although finite, very small, typically of order $\lesssim 1$ eV, and can be made consistent with present experimental bounds and cosmological arguments.

It was shown[5] recently that cosmological excess of baryons over antibaryons in the present universe may be explained by baryon- and CP-nonconserving processes in the very early universe. The key number of baryon to photon ratio N_B/N_γ of order 10^{-9} may be obtained[6] by either heavy ($\approx 10^{15}$ GeV) gauge or heavy Higgs boson decay when they go out of thermal equilibrium. It was, however, necessary[6] for this purpose to have more than a minimal set of Higgs bosons, for instance, more than two quintets $(\underline{5})$ of Higgs bosons in the SU(5) model.

Here we wish to point out that an alternative mechanism of baryon production via decay of heavy Majorana particles is possible without introducing additional Higgs bosons besides the minimal set. Processes that may contribute to the baryon excess in this class of models are three-quark decays,

$$N \to qqq, \; \bar{q}\bar{q}\bar{q}. \tag{3}$$

Two decay modes here with different baryon numbers can occur because Majorana fields N do not have a definite particle number. An example of diagrams that contribute in leading order to the baryon asymmetry ΔB is shown in Fig. 1, and in the notation of Eq. (1) gives an asymmetry proportional to an imaginary part of couplings,

$$\Delta B = \pi^{-3} \textstyle\sum_{ij} [\mathrm{tr}\lambda\lambda^\dagger \bar{M}_i]^{-1} \,\mathrm{Im}\,\mathrm{tr}[\lambda^T \lambda^* \bar{M}_i \lambda^\dagger f f^\dagger \lambda \bar{M}_j] F(\bar{M}_i, \bar{M}_j), \tag{4}$$

where $\bar{M}_i$ is the ith diagonal element of Majorana mass matrix divided by the mass of the colored Higgs boson H in Fig. 1. We have used the fact that a main two-body decay mode of N is $O(\lambda)$ in amplitude. An important point is the necessity of the discontinuity as shown in the dot-dashed line of Fig. 1, which is required to give different decay rates between the two modes of Eq. (3). Existence of the discontinuity means that $F(\bar{M}_i, \bar{M}_j)$ is not symmetric under interchange of $\bar{M}_i$ and $\bar{M}_j$ because one is larger than the other. Hence, Eq. (4) does not vanish unless CP-noncon-

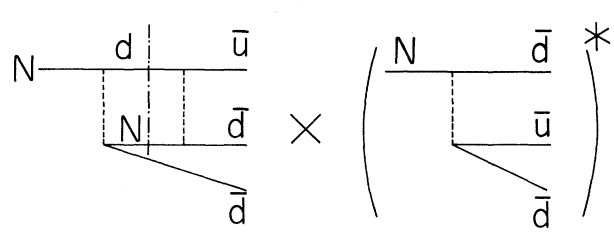

FIG. 1. Example of interference diagrams that contributes to the baryon asymmetry via $N \to \bar{u}\bar{d}\bar{d}$. Dotted lines represent colored Higgs bosons and the dot-dashed line a discontinuity to be taken.

serving complex parameters conspire to cancel their phases in Eq. (4).

To get a rough idea of the magnitude of Eq. (4) we note that $F(\bar{M}_i, \bar{M}_j)$ is of order unity when the arguments are of order unity. In this case,

$$\Delta B = O[\langle \lambda\lambda^\dagger ff^\dagger \rangle / \pi^3] \tag{5}$$

unless CP-invariance violation is much suppressed from unknown reasons. Here $\langle \cdots \rangle$ means an average over various couplings. As shown by the works of Refs. 5 and 6, this asummetry ΔB is related to the baryon to photon ratio N_B/N_γ by a calculable factor of order 10^{-2} to 10^{-3}. It is therefore necessary that

$$\langle \lambda\lambda^\dagger ff^\dagger \rangle / \pi^3 \approx 10^{-6} \tag{6}$$

for this particular mechanism of baryon production to work. Typical magnitudes of Yukawa couplings are expected to be of order

$$\langle ff^\dagger \rangle / 4\pi \approx \langle hh^\dagger \rangle / 4\pi \approx 10^{-3}. \tag{7}$$

Hence, Eq. (6) is viable if λ is of a similar magnitude to f or h. Unlike the case considered above, if all $|\bar{M}_i| \ll 1$, then ΔB given by Eq. (4) is further suppressed by factors like $|\bar{M}_i|^2$. We may then conclude that with typical magnitudes of Yukawa couplings, at least some of Majorana particles must be roughly as heavy as the colored Higgs boson in order to obtain a reasonable, cosmological baryon asymmetry.

Kinematical consideration[5] that arises from the nonequilibrium nature of the Majorana decay in the above case leads to a mild constraint on the colored Higgs mass $m_{\rm H}$,

$$m_{\rm H} \gtrsim O[\{\langle \lambda\lambda^\dagger (ff^\dagger + hh^\dagger) \rangle\}^{1/4} (\langle M \rangle / m_*)^{3/3} m_* / 4\pi], \tag{8}$$

with the Planck mass m_*. The fact that this gives an upper bound to an average of the Majorana mass $\langle M \rangle$ is essentially due to the 4-Fermi nature of baryon-nonconserving decay $N \to qqq$.

Arguments for the cosmological baryon production thus suggest that

$$\langle M \rangle \approx m_{\rm H} \gtrsim 10^{11}\ \text{GeV}, \quad \langle \lambda\lambda^\dagger \rangle / 4\pi \approx 10^{-3} \tag{9}$$

if ordinary Yukawa couplings are of the order of Eq. (7). This constraint on the Higgs mass $m_{\rm H}$ is roughly the same as the bound given by proton decay. With these huge masses of Majorana leptons, flavor mixings among charged leptons that are induced in higher orders of perturbation are very much suppressed either by small neutrino masses or by large Majorana masses. What about neutrino oscillations? The oscillation length with a neutrino energy $E \gg |m_\nu|$ is given by[7]

$$L = 4\pi E / \langle (\Delta m)^2 \rangle, \tag{10}$$

where $\langle (\Delta m)^2 \rangle$ is an average of differences of neutrino masses squared. It is perhaps natural to expect that

$$\langle (\Delta m)^2 \rangle \approx \langle m_\nu^2 \rangle \approx [v^2 \langle \lambda^2 \rangle / \langle M \rangle]^2.$$

From the kinematical constraint (8) with $m_{\rm H} \approx \langle M \rangle$, the right-hand side of this equation is bounded from above and

$$\langle (\Delta m)^2 \rangle \lesssim O((4\pi)^8 v^4 / (\langle ff^\dagger \rangle^2 m_*^2)), \tag{11}$$

which is numerically of order 1 eV^2 with Eq. (7). We stress that this limit (11) is almost independent of the unknown magnitudes, λ and M. This size of oscillation length is just above the lower limit[7] that present reactor and accelerator experiments impose. On the other hand, the solar-neutrino "puzzle"[4] may already be an indication of neutrino oscillation with an oscillation length $\lesssim 10^{13}$ cm. If we use this as a positive evidence of neutrino oscillation, we may set a rough upper bound of the Majorana mass, $\langle M \rangle$ which is of order $\lesssim 10^{17}$ GeV, assuming the magnitude of λ given by Eq. (9). In view of a wide gap between this mass and the one given by Eq. (9), it would be extremely helpful if limits by terrestrial experiments were further improved.

We have also considered within our model other possibilities for mechanisms of cosmological baryon production. For instance, when the colored Higgs boson H is less massive than a Majorana particle N, the two-body decays given by $N \to H\bar{q}$ and $N \to \bar{H}q$ may become a dominant process of baryon production. In this case the asymmetry ΔB is of order $\langle \lambda\lambda^\dagger \rangle / 16\pi$ and the kinematical constraint leads to

$$M \gtrsim 0.01 \lambda\lambda^\dagger m_* / 4\pi. \tag{12}$$

This gives a rough upper limit of neutrino masses of order 10^{-3} eV. Interestingly enough, this value is independent of almost all the details of coupling strengths and corresponds to an oscillation length, roughly of order 10^6 km for a neutrino energy of 1 GeV.

Another possible implication of the Majorana particle considered here is violation of the selection rule[8] $B-L$ in nucleon decay. This violation occurs in one-loop diagrams for $dd\to\bar{d}e$ in which the heavy Majorana as well as the colored Higgs are circulated. But its magnitude appears too small compared with that of the main nucleon decay to have any chance of detection.

In summary, we have demonstrated a possibility that both the cosmological baryon excess and the solar-neutrino "puzzle" may be explained by heavy Majorana particles of masses $\gtrsim 10^{10}$ GeV. Our model is perhaps too simple and our numerical estimates are admittedly very crude, which we feel to be justified by great uncertainties surrounding the grand unified theories. Forthcoming experiments on nucleon decay and neutrino oscillation are expected to shed more light on the questions raised in this paper.

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ERRATUM

DIRECT OBSERVATION OF FINE STRUCTURE OF MAGNETIC DOMAIN WALLS BY ELECTRON HOLOGRAPHY. Akira Tonomura, Tsuyoshi Matsuda, Junji Endo, Tatsuo Arii, and Kazuhiro Mihama [Phys. Rev. Lett. 44, 1430 (1980)].

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