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MATTER-ANTIMATTER OSCILLATIONS AND NEUTRINO MASS*

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Goran Senjanovic

Physics Department
Brookhaven National Laboratory
Upton, NY 11973

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MATTER - ANTIMATTER OSCILLATIONS AND NEUTRINO MASS

Goran Senjanović

Brookhaven National Laboratory, Upton, New York 11973, U.S.A.

ABSTRACT

A discussion of neutron-antineutron ($n-\bar{n}$) and hydrogen-antihydrogen ($H-\bar{H}$) transitions is presented. An $SU(2)_L \times U(1) \times SU(3)_C$ model with spontaneously broken global B-L symmetry is shown to predict the interesting connection between oscillation times $T_{n-\bar{n}}$, $T_{H-\bar{H}}$, neutrino mass and the mass of a doubly charged Higgs scalar. A case of B-L as a gauge symmetry is discussed in the context of $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_C$ gauge model, with the emphasis on matter oscillations. Finally, an analysis of Higgs mass scales in GUTS and their impact on such processes is offered.

I. INTRODUCTION

Now that the idea of grand unification¹ has become so well established, one of the main questions that it poses is which kind of baryon number (B) and lepton number (L) violating processes is to be observed in the near future. A few years ago Weinberg and Wilczek and Zee have offered² a model independent way of analyzing that question based on the short distance expansion of the leading effective baryon number violating operators. If an assumption of a single mass scale beyond M_W is made, the leading operators are those of the lowest dimension and they take the necessary form which conserves B-L

$$O = \frac{1}{M_X^2} (qqq\ell) \quad (1)$$

where q stands for quarks and ℓ for leptons and M_X is the large scale beyond M_W . As is well known, these operators correspond to the $\Delta B = 1$, B-L conserving nucleon decay $n \rightarrow \pi^- + e^+$, $p \rightarrow \pi^0 + e^+$ and the limits on nucleon stability³ $\tau_p \gtrsim 10^{31}$ yr imply $M_X \gtrsim 10^{14}-10^{15}$ GeV. Using Georgi, Quinn and Weinberg program⁴ in GUTS with a single scale above M_W , such as $SU(5)$, M_X is predicted from unification constraints to be between $10^{14}-10^{15}$ GeV; a spectacular prediction to be tested in the near future.

If the above is true, we will encounter the unprecedented situation in the history of physics: an enormous "desert" in energies between $100-10^{14}$ GeV is predicted in which no new physics is to be found! No matter how exciting the prediction of proton decay is, this would certainly make the future of particle physics rather grim. Various scenarios that suggest filling the desert with one or more oasis, such as technicolor⁵ or supersymmetric GUTS⁶ have come forward, however, yet without full success. A more modest approach would be to search for an oasis within the context of conventional grand unification, an approach that I will follow here in my talk and an approach more suitable to the character of this conference.

To understand better the emergence of a desert and the relevance of B-L symmetry, we will start with the discussion of B and L nonconservation in the standard $SU(2)_L \times U(1) \times SU(3)_C$ model. Although perturbatively the above quantum numbers are absolutely conserved, 't Hooft has demonstrated that the nonperturbative, instanton induced effects, lead to the violation of B and L, but still conserve B-L global symmetry. Again, B-L emerges as the relevant symmetry! It appears then natural to study the question of an oasis in the desert through the properties of B-L operator and this study comprises the main part of this paper.

We start first with the discussion of B-L conservation in the standard model, as presented in Section II. If, as it is natural in gauge theories, B-L is broken spontaneously, an interesting possibility of hopefully observable $n-\bar{n}$ oscillations (which violate B-L) and $H-\bar{H}$ oscillations (which conserve B-L) appear in the context of the extended $SU(2)_L \times U(1)$ model. The relevant transition times, although not predicted in this model, become related to the neutrino mass and the mass of the doubly charged Higgs scalar. This model is easily extended to simple GUTS, such as $SU(5)$. Next, in Section III, a different possibility, in which B-L becomes a part of a gauge electroweak group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ is briefly reviewed. These left-right symmetric models have been extensively analyzed in the past and we cover them here from the point of view of B-L breaking, with the emphasis on $n-\bar{n}$ and $H-\bar{H}$ transitions. In Section IV we discuss the question of oscillation times expected in simple GUTS. This issue is intimately tied up to the question of neutral Higgs mass scales in GUTS. It is agreed that a proper treatment of Higgs boson masses renders these processes unobservable, at least in simple grand unified models. Finally, in Section V a summary of this paper is offered.

§II. THE STANDARD MODEL AND GLOBAL B-L SYMMETRY BREAKING

Let us first briefly recall the main features of the minimal $SU(2)_L \times U(1) \times SU(3)_C$ model which will also be useful in setting our notations,

Fermions

$$\psi_L = \begin{pmatrix} \nu \\ e \end{pmatrix}_L \quad (2, -1, 1_C) \quad ; \quad e_R \quad (1, -2, 1_C) \quad (2.1)$$

$$Q_L = \begin{pmatrix} u \\ d \end{pmatrix}_L \quad (2, 1/3, 3_C) \quad ; \quad u_R \quad (1, 4/3, 3_C), \quad d_R \quad (1, -2/3, 3_C)$$

Higgs

$$\phi = \begin{pmatrix} \phi^+ \\ \phi_0 \end{pmatrix} \quad (1, 2, 1_C) \quad (2.2)$$

The most general Yukawa interaction is

$$L_Y = h_u \bar{Q}_L i\tau_2 \phi^* u_R + h_d \bar{Q}_L \phi d_R + h_e \bar{Q}_L \phi e_R + \text{h.c.} \quad (2.3)$$

As encouraged by experiment, ν_R is assumed not to exist and neutrino is massless to all orders in perturbation theory. It is easy to see that (2.3) and the rest of the Lagrangian is invariant under $(U1)_B$ and $U(1)_L$ global symmetries, with $B\phi = L\phi = 0$. A minimal realistic theory of quarks and leptons appears to naturally conserve baryon and lepton numbers.

However, the above is not the whole story. As is well known, if a global current has an axial anomaly, the current conservation does not hold true, i.e. one gets⁸

$$\partial^\mu J_\mu = \frac{g^2}{16\pi^2} A F^{\mu\nu} F_{\mu\nu}^d \quad (2.4)$$

where $F_{\mu\nu}$ is an Abelian or non-Abelian field strength and $F_{\mu\nu}^d = \epsilon_{\mu\nu\alpha\beta} F^{\alpha\beta}$ and A is the anomaly coefficient. From (2.4) it is easy to see that the charge $Q = \int d^3x J^0$ is not conserved in that case, i.e.

$$\frac{d}{dt} Q = \frac{g^2}{16\pi^2} A \int d^3x F^{\mu\nu} F_{\mu\nu}^d \quad (2.5)$$

Although the integrand on the right hand side is the total derivative, in the presence of instantons the integral does not vanish. Now, the anomaly coefficient $A \propto \text{Tr} Q T_i T_j$ and so from, say

$$\begin{aligned} \text{Tr } B T_{3L}^2 &= \frac{1}{3} \frac{1}{2} \cdot 3 = \frac{1}{2} \\ \text{Tr } L T_{3L}^2 &= 1 \cdot \frac{1}{2} = \frac{1}{2} \end{aligned} \quad (2.6)$$

B and L are not anomaly free and so are not conserved by the nonperturbative effects⁹. Notice, however that the $B-L$ combination is anomaly free. The breakdown of B and L is in this case only of academic interest, since the nonperturbative effects are enormously suppressed by $e^{-4\pi\alpha/\sin^2\theta_W}$, where α is the electromagnetic coupling. Nevertheless, the special character of $B-L$ is again seen.

Now, to get around the desert picture, $B-L$ must be broken. Since the explicit breaking lacks any predictive power, in what follows we shall pursue the idea of the spontaneous breakdown of $B-L$ symmetry⁹.

Extended $SU(2)_L \times U(1) \times SU(3)_c$ model and the spontaneous breakdown of global $B-L$ symmetry: This model has been developed in Ref. 9 and the phenomenological implications presented here were worked out by Mohapatra and myself⁹. Our main idea was to incorporate the notion of spontaneously broken global $B-L$ symmetry into a study of

neutron (n - $\bar{n}$) and hydrogen (H - $\bar{H}$) oscillations; the result of which, as we shall see is the connection between these and the neutrino Majorana mass.

Let us start by describing the model. Besides the usual particles given in (2.1) and (2.2), in order to break B-L spontaneously, one is forced into extending the Higgs sector with the following $SU(2)_L$ triplets⁹.

$$\Delta_\ell(3, 2, 1_c) ; \Delta_q(3, -\frac{2}{3}, \bar{6}_c) \quad (2.7)$$

The existence of Δ_ℓ , suggested in Ref. 10, is sufficient to break B-L; but Δ_q enables the preservation of quark-lepton symmetry of the standard model, and furthermore will lead to $n-\bar{n}$ and $H-\bar{H}$ transitions.

We assign $B_{\Delta_q} = -2/3$, $L_{\Delta_q} = 0$, and $B_{\Delta_\ell} = -2$, and demand the full Lagrangian to possess global $U(1)_{B-L}$ symmetry. The most general Yukawa interaction and Higgs potential are then given by ($\Delta \equiv 1/\sqrt{2} \vec{\tau} \cdot \vec{\Delta}$)

$$L_Y = h_u \bar{Q}_L (i\tau_2 \phi^*) u_R + h_d \bar{Q}_L \phi^d_R + h_e \bar{\psi}_L \phi^e_R \\ + f_q Q_L^T C^{-1} \tau_2 \Delta_q Q_L + f_\ell \psi_L^T C^{-1} \tau_2 \Delta_\ell \psi_L + h.c. \quad (2.8)$$

and

$$V = -\mu_\phi^2 \phi^\dagger \phi + \lambda_\phi (\phi^\dagger \phi)^2 \\ - \mu_\ell^2 \text{Tr} \Delta_\ell^\dagger \Delta_\ell + \lambda_\ell (\text{Tr} \Delta_\ell^\dagger \Delta_\ell)^2 + \lambda'_\ell \text{Tr} \Delta_\ell^\dagger \Delta_\ell \Delta_\ell^\dagger \Delta_\ell \\ - \mu_q^2 \text{Tr} \Delta_q^\dagger \Delta_q + \lambda_q (\text{Tr} \Delta_q^\dagger \Delta_q)^2 + \lambda'_q \text{Tr} \Delta_q^\dagger \Delta_q \Delta_q^\dagger \Delta_q \\ + \alpha \text{Tr} \Delta_\ell^\dagger \Delta_\ell \text{Tr} \Delta_q^\dagger \Delta_q + \alpha' \text{Tr} \Delta_q^\dagger \Delta_q \Delta_\ell^\dagger \Delta_\ell \\ + \phi^\dagger \phi (\beta_\ell \text{Tr} \Delta_\ell^\dagger \Delta_\ell + \beta_q \text{Tr} \Delta_q^\dagger \Delta_q) + \phi^\dagger (\gamma_\ell \Delta_\ell^\dagger \Delta_\ell + \gamma_q \Delta_q^\dagger \Delta_q) \phi \\ \lambda \epsilon_{ikm} \epsilon_{jln} (\text{Tr} \Delta_q^{ij} \Delta_q^{kl}) (\text{Tr} \Delta_q^{mn} \Delta_\ell) + h.c. \quad (2.9)$$

where we have explicitly included the only nontrivial colour indices in the last term of (2.9). That term is very important: in its absence the Lagrangian admits two separate $U(1)$ symmetries corresponding to baryon and lepton numbers; however, $\lambda \neq 0$ leaves only $U(1)_{B-L}$ broken.

Incidentally, the theory as a by-product possesses a discrete symmetry D_B , with all the fields having the transformation property: (field) $\rightarrow e^{i\pi B}$ (field). In particular, under D_B , $p \rightarrow -p$, $n \rightarrow -n$, but $e \rightarrow e$ and mesons $\rightarrow$ mesons. This symmetry, as we shall see below, plays a crucial role in guaranteeing the stability of the proton.

The potential in Eq. (2.9) is minimized for the following values of the fields, which break the global B-L symmetry¹⁰:

$$\langle \phi \rangle = \begin{pmatrix} 0 \\ v/\sqrt{2} \end{pmatrix}, \quad \Delta_\ell = \begin{pmatrix} 0 & 0 \\ \kappa/\sqrt{2} & 0 \end{pmatrix} \quad (2.10)$$

leading to the usual expressions for quark and electron masses and a Majorana mass for the neutrino

$$m_\nu = \sqrt{2} f_e \kappa \quad (2.11)$$

The leptonic aspect of this model has been discussed at length before¹⁰⁻¹². The spontaneous breakdown of global B-L symmetry leads to an existence of a zero mass Goldstone boson, the Majorana^{13,10}, whose presence requires re-doing some of the usual weak interaction phenomenology. The only constraints of importance for us is $f_e < 7 \times 10^{-3}$ from π and K meson decays¹² and $\kappa < 100$ keV from the astrophysical analysis of energy loss from red giants^{11,14}. This results in $m_{\nu_\ell} < 700$ eV; obviously not a very useful bound.

Matter-antimatter transitions: We now turn our attention to $\Delta B = 2$, $\Delta L = 0$ $n\bar{n}$ transitions¹⁵ and $\Delta B = \Delta L = 2$ $H\bar{H}$ transitions¹⁶. Notice that, since a discrete symmetry D_B remains unbroken, proton is absolutely stable. That means that we have no reason to assume the existence of new large mass scales in the theory; in what follows we shall make the most natural assumption of a single mass scale, i.e. that of weak interactions ~ 100 GeV.

$H\bar{H}$ oscillations: This process proceeds in the manner shown in Fig. 1. Its strength is easily estimated to be

$$G_{H\bar{H}} = \frac{\lambda f_q^3 f_e}{m_{\Delta q}^6 m_{\Delta^{++}}^2} \quad (2.12)$$

This in turn leads to an effective hydrogen-antihydrogen mixing of the strength

$$H_{\text{eff}} = G_{H\bar{H}} |\psi_N(0)|^4 (p_L^T C^{-1} p_L)(e_L^T C^{-1} e_L) + \text{h.c.} \quad (2.13)$$

where $|\psi_N(0)|^4$ is the quark wave function inside the nucleus, and e and p are electron and proton fields. From Ref. 16 we can then estimate the mixing time for $H\bar{H}$ oscillation

$$T_{H\bar{H}}^{-1} = G_{H\bar{H}} |\psi_N(0)|^4 \frac{(m_e \alpha)^3}{\pi} \quad (2.14)$$

In what follows we will take $m_{\Delta q} \sim 100$ GeV and $\lambda \sim 1$ as natural

values. Since $f_e < 7 \times 10^{-3}$, $m_{\Delta^{++}} > 15 \text{ GeV}$,¹¹ by assuming $f_q < 10^{-3}$ and taking $^{17} |\psi_N(0)|^4 \Delta^{++} \approx 10^{-5} (\text{GeV})^6$ we get from (2.12) and (2.14)

$$T_{H-\bar{H}} \geq 3 \times 10^{15} \text{ yr} \quad (2.15)$$

Fig. 1 Example of a quark diagram involving $H \rightarrow \bar{H}$ process. Δ_{uu} Δ_{dd} are charge $-4/3$ and $+2/3$ members of the Δ_q triplet, and Δ^{++} is the doubly charged Higgs scalar from Δ_ℓ .

We wish to add that, as is usual with $\Delta B \neq 0$ processes, our prediction depends sensitively on the choice of Higgs particle masses; however, our estimate is fairly reasonable. We eliminate a large degree of uncertainty when we make comparison of $T_{H-\bar{H}}$ with (possibly) observable $T_{n-\bar{n}}$.

$n-\bar{n}$ oscillations: It is easy to estimate $T_{n-\bar{n}}$ from Fig. 2, in exactly the same manner as for $H-\bar{H}$ transitions

Fig. 2 A basic quark diagram depicting $n-\bar{n}$ transitions.

$$T_{n-\bar{n}}^{-1} = \frac{\lambda f_q^3}{m_{\Delta_q}^6} \frac{m_\nu}{\sqrt{2} f_e} |\psi_N(0)|^4 \quad (2.16)$$

where we have used (2.11). For the same values of parameters as before, we obtain

$$T_{n-\bar{n}} \approx 10^{22} m_\nu^{-1} \quad (2.17)$$

The experimental limit¹⁸ $T_{n-\bar{n}} > 10^7-10^8$ sec requires then $m_\nu < 10^{-1}$ eV (of course, the smaller couplings and/or larger value of m_{Δ_q} would suppress the above process and allow for heavier neutrinos).

From (2.12), (2.14) and (2.16) we arrive at our main prediction

$$\frac{T_{H-\bar{H}}}{T_{n-\bar{n}}} = \frac{\pi m_{\Delta^{++}}^2 m_\nu}{f_e^2 (\alpha m_e)^3} \quad (2.18)$$

which is free from most of the uncertainties that plagued $T_{n-\bar{n}}$ and $T_{H-\bar{H}}$. For example, if we choose $m_{\Delta^{++}} \approx 10-100$ GeV, we expect

$$\frac{T_{H-\bar{H}}}{T_{n-\bar{n}}} \approx 5 \times 10^{20} \frac{m_\nu (\text{GeV})}{m_e} \quad (2.19)$$

For $T_{n-\bar{n}} \approx 1$ yr, $m_\nu \approx (1-10)$ eV, $T_{H-\bar{H}} \approx 10^{13}$ yr, which is close to the lower bound inferred from astrophysics by Feinberg et al.¹⁶. The important point to note is that acceptable and interesting baryon non-conserving processes can be predicted in theories without any mass scale beyond $M \approx M_W$.

As far as laboratory experimental searches for $H-\bar{H}$ mixing is concerned, long mixing times would appear to make it inaccessible to observation in existing set ups. However, the processes such as $p + n \rightarrow e^+ + \nu$ and $n \rightarrow p + e^+ + \nu$, and especially double proton decay described below, are free from the suppression due to atomic wave functions and so could be searched for in the same set ups that look for proton decays.

Double proton decay: $p + p \rightarrow e^+ + e^+$: The process proceeds through the same diagram in Fig. 1 that describes $H-\bar{H}$ transition and can be characterized by the same effective strength $G_{H-\bar{H}}$ defined in (2.12). For the same values of the parameters as before, we estimate in our model $G_{H-\bar{H}} |\psi_N(0)|^4 < 10^{-29} (\text{GeV})^{-2}$, or in other words $\tau_{pp} > 10^{31}$ yr, close to the experimental bound.

For values of $T_{H-\bar{H}}$ that we predict, there also exists the possibility of observing $H-\bar{H}$ mixing effects in astrophysics by looking for energetic interstellar γ -rays. As for other distinguishable features of this model, we note that:

- i) doubly charged Higgs bosons (with $m_{\Delta^{++}} = 15$ GeV or so) should be visible in existing high energy machines in the not too distant future;
- ii) electron neutrino should weight about 10^{-1} eV, if both $n-\bar{n}$ and $H-\bar{H}$ processes are to have detectable rates;
- iii) if there exist direct $\Delta S = 2$ tree level amplitude through the exchange of a diquark Δ_{dd} , with the strength $f_d f_s / m^2 \Delta_q$. For $f_d < 10^{-3}$ and $m_{\Delta_q} \approx 100$ GeV, we need $f_s / f_d < 10^{-3}$ in order to guarantee the smallness of $K_L - K_S$ mass difference. This would predict the large suppression of possible $\Lambda - \bar{\Lambda}$ transition, compared to $n - \bar{n}$ process

In closing we should mention that this model provides the simplest framework which realizes the general possibility of mutually

related $T_{n-\bar{n}}$ and $T_{H-\bar{H}}$ when only left-handed operators contribute. The connection is necessary since we work in the context of purely left-handed, standard model.

SU(5) model with spontaneously broken global B-L symmetry: We have shown in the previous section how a fairly simple extension of the Higgs sector in the standard model can lead to mutually connected and possibly observable $n-\bar{n}$ and $H-\bar{H}$ oscillations. Here, we shall discuss its grand unified version.

Gauge group: $G = SU(5) \times U(1)_G$ with $U(1)_G$ a global symmetry specified below; we shall call its generator G .

Fermions: As in the minimal SU(5) model¹⁹, we have

$$\underline{5} = F_R = \begin{pmatrix} d_1 \\ d_2 \\ d_3 \\ e^+ \\ \nu^c \end{pmatrix}_R$$

$$\underline{10} = T_L = \begin{pmatrix} 0 & u_3^c & -u_2^c & -u_1 & -d_1 \\ -u_3^c & 0 & u_1^c & -u_2 & -d_2 \\ u_2^c & -u_1^c & 0 & -u_3 & -d_3 \\ u_1 & u_2 & u_3 & 0 & e^+ \\ d_1 & d_2 & d_3 & -e^+ & 0 \end{pmatrix}_L$$

Higgs sector: Apart from $\underline{5}$ dimensional representation H and $\underline{24}$ dimensional in representation Σ of the minimal model, we include the following additional Higgs multiplets to implement our idea of spontaneous B-L symmetry breakdown²⁰

$$\begin{aligned} \underline{15} & : \Delta_{pp} \\ \underline{50} & : \chi_{rst}^{pq} \end{aligned} \quad (2.21)$$

where Δ_{pq} is symmetric: $\Delta_{pq} = \Delta_{qp}$ and χ_{rst}^{pq} is totally antisymmetric

upper and lower indices, respectively and is traceless: $\chi_{pst}^{pq} = 0$.

The most general Yukawa interaction consistent with SU(5) gauge symmetry is

$$\begin{aligned} L_Y = & h_d F_R^P H^q T_{Lpq} \\ & + h_u \epsilon^{pqrst} T_{Lpq}^T H_r C^{-1} T_{Lst} + h_\nu F_{Rp}^T C^{-1} \Delta^{pq} F_{Rq} \\ & + f \epsilon^{pqrst} T_{Lpp_1}^T C^{-1} \chi_{rst}^{p_1 q_1} T_{Lqq_1} + \text{h.c.} \end{aligned} \quad (2.22)$$

It is easy to see that the above interaction is invariant under a $U(1)_G$ global symmetry $U = e^{i\theta G}$, where

$$\begin{aligned} G T_L &= T_L, \quad G F_R = 3 F_R \\ G H &= -2 H, \quad G \Delta = G \chi = -2 \chi \end{aligned} \quad (2.23)$$

all other fields of the theory having no G charge. The invariance of the rest of the Lagrangian under this symmetry is not automatic. It requires omitting various renormalizable couplings such as $H \Delta^\dagger H$, $\Sigma \Delta^\dagger \Sigma$, etc. Therefore, we demand $U(1)_G$ to be a symmetry of the full Lagrangian.

Now, $\langle \Sigma \rangle \approx M_X$ which is responsible for breaking of SU(5) down to $SU(2)_L \times U(1) \times SU(3)_C$ will not break $U(1)_G$. On the other hand, $\langle H \rangle \approx M_W$ which is responsible for the second stage of the symmetry breaking, breaks $U(1)_G$, since $G \langle H \rangle = -2 \langle H \rangle$ but preserves $(Y + G/2) \cdot \langle H \rangle = \langle H \rangle$. A simple computation shows¹⁸ that this is nothing else but B-L symmetry when it acts on fermions

$$B-L = \frac{2}{5} \left(Y + \frac{G}{2} \right) \quad (2.24)$$

The potential of the theory can be chosen in the conventional manner to produce a minimum at $\langle \Delta_{55} \rangle \neq 0$, which leads to the spontaneous breakdown of B-L symmetry manifest through a Majorana mass for the neutrino

$$m_\nu = h_\nu \langle \Delta \rangle \quad (2.25)$$

The associated Goldstone boson, the majoron is now given by

$$M = \frac{\langle H \rangle \operatorname{Im} \Delta_{55} - \langle \Delta \rangle \operatorname{Im} H_5}{(\langle H \rangle^2 + \langle \Delta \rangle^2)^{1/2}} \quad (2.26)$$

and its properties are essentially the same as in the $SU(2)_L \times U(1)$ model.

We now turn our attention to matter oscillations. These processes, as we show below, are mediated by particles in Δ and χ multiplets. Let us then first display the $SU(2)_L \times U(1) \times SU(3)_C$ representation contents of Δ and χ

$$\begin{aligned} & (3, 2, 1_c) + (1, -4/3, 6_c) + (2, 1/3, 3_c) \\ \Delta = & \Delta_{ab} = \Delta_1 \quad \Delta_{ij} = \Delta_2 \quad \Delta_{ia} = \Delta_3 \end{aligned} \quad (2.27)$$

where $a, b, \dots$ are $SU(2)_L$ indices and $i, j, k, \dots$ are $SU(3)_C$ indices. Similarly, for χ we get

$$\begin{aligned} & (1, -4, 1_c) + (1, -2/3, 3_c) + (2, -7/3, 3_c) \\ \chi = & \chi_{ijk}^{ab} = \chi_1 \quad \chi_{cdi}^{ab} = \chi_2 \quad \chi_{cij}^{ab} = \chi_3 \\ & (1, 8/3, 6_c) + (3, -2/3, 6_c) + (2, 1, 8_c) \\ + & \chi_{kab}^{ij} = \chi_4 \quad \chi_{bjk}^{ai} = \chi_5 \quad \chi_{jbc}^{ia} = \chi_6 \end{aligned} \quad (2.28)$$

Before we discuss any physical consequences due to these Higgs scalars, we would have to know their mass scales. If one assumes minimal fine tuning, .e. one makes no more fine tuning of the parameters beyond the one needed for the gauge hierarchy $M_W/M_X \approx 10^{-12}$, then the survival hypothesis is operative: only in those particles which have to be light will remain light (see Section IV). In this case it is the $SU(2)_L$ doublet (H_4, H_5) from H and the $SU(2)_L$ triplet Δ_1 from Δ . Namely (H_4, H_5) contains Goldstone bosons longitudinal components of $W^\pm$ and Z) and the physical scalar η ; whereas Δ_1 contains the majoron to prevent the rest of the multiplet from becoming superheavy. All other Higgs scalars would however become heavy. This of course, would imply the suppression of all other processes but proton decay and we somewhat uneasily give it up. We could imagine that some symmetry (maybe supersymmetry) could eventually keep other particles light, but for the moment we just assume it (it is always technically possible to do it by adjusting the parameters). Still, however, χ_2 has to be superheavy because it mediates proton decay and so $m_{\chi_2} = M_X$. As far as other particles are concerned, we will keep them light and in particular we assume that χ_1, χ_4, χ_5 and Δ_1 submultiplets have masses of order 100 GeV.

It is worth noting that χ_1 and Δ_2 are doubly charged Higgs bosons, where χ_1 couples to right-handed leptons and Δ_1 couples to the left-handed ones and is the analogue of the Δ_L of Section III. Since this is somewhat conceptual, I will not pursue it here. Suffice it to say that both $n\bar{n}$ and $H\bar{H}$ transitions proceed in the manner described in this section. Furthermore, the model can be easily made consistent with constraints as $\sin^2\theta_W$ and T_R . For the details, the interested reader should consider a paper²² of Mohapatra and myself which discusses these questions at length.

III. B-L AS A GAUGE GENERATOR: LEFT-RIGHT SYMMETRY

In the previous section we have reviewed the possibility of B-L being a global symmetry, subsequently spontaneously broken by a Higgs mechanism. This led to a plethora of interesting new phenomena, in particular Majorana neutrino mass, existence of Goldstone bosons (Majorons) and the connection between m_ν and $T_{n\bar{n}}$, $T_{H\bar{H}}$. The other, equally interesting possibility is that B-L is a gauge symmetry, again spontaneously broken. This actually happens naturally in a class of gauge theories, characterized by left-right symmetry²³. We offer a brief review of such models here, with the emphasis on B-L breaking.

The gauge groups which incorporate this idea is $SU(2)_L \times SU(2)_R \times U(1)_{B-L} \times SU(3)_C$ with the fermionic assignment²³

$$Q_L = \begin{pmatrix} u \\ d \end{pmatrix}_L \quad (2, 1, \frac{1}{3}, 3_C) ; \quad Q_R = \begin{pmatrix} u \\ d \end{pmatrix}_R \quad (1, 2, \frac{1}{3}, 3_C)$$

$$\psi_L = \begin{pmatrix} \nu \\ e \end{pmatrix}_L \quad (2, 1, -1, 1_C) ; \quad \psi_R = \begin{pmatrix} \nu \\ e \end{pmatrix}_R \quad (1, 2, -1, 1_C) \quad (3.1)$$

which is completely left-right symmetric as is the rest of the theory. The electromagnetic charge in this model is seen to be²⁴ from (3.1)

$$Q_{em} = I_{3L} + I_{3R} + \frac{B-L}{2} \quad (3.2)$$

so that the $U(1)$ generator is physical: it is B-L. The above relation indicates the clear connection between the breaking of parity (I_{3R}) and B-L. This was realized in the work of Mohapatra and the author²⁵, through the connection of the smallness of neutrino Majorana mass and the maximality of parity violation in weak interactions. In order to describe it we must introduce the minimal Higgs sector which incorporates²⁵ the breaking of parity and B-L

$$\Delta_L^{\frac{1}{2}} (3, 1, 2, 1_C) ; \quad \Delta_R^{\frac{1}{2}} (1, 3, 2, 1_C) ; \quad \phi (2, 2, 0, 1_C) \quad (3.3)$$

Subsequent symmetry breaking

$$\langle \Delta_L^f \rangle = 0, \quad \langle \Delta_R^f \rangle = g M_R, \quad \langle \phi \rangle = g M_W \quad (3.4)$$

in the symbolic notation gives a realistic theory with broken parity, through $M_R > M_W$. What is more interesting to us, from the Yukawa interaction for leptons

$$L_Y = f_e (\psi_L^T \tau_2 C \Delta_L \psi_L + \psi_R^T \tau_2 C \Delta_R \psi_R) + h \bar{\psi}_L \phi \psi_R \quad (3.5)$$

where

$$\Delta_{L,R} \equiv \frac{1}{\sqrt{2}} \vec{\tau} \cdot \vec{\Delta}_{L,R},$$

right-handed neutrino gets a large Majorana mass

$$m_{\nu_R} = f M_R \quad (3.6)$$

and in turn the left-handed neutrino becomes naturally small

$$m_{\nu_L} = \frac{h^2 \langle \phi \rangle^2}{m_{\nu_R}} \quad (3.7)$$

In the V-A limit of the theory, (i.e. $M_R \rightarrow \infty$), $m_{\nu_R} \rightarrow \infty$ and $m_{\nu_L} \rightarrow 0$; hence, the smallness of neutrino mass gets tied up to the maximality of observed parity violation in weak interactions.

An interesting extension of the minimal model has been suggested by Marshak and Mohapatra²⁶, by making the model completely quark-lepton symmetric. They add, besides leptonic triplets Δ^f , quark triplets which are sextets of colour

$$\Delta_L^q (3, 1, -\frac{2}{3}, \bar{6}_c) ; \quad \Delta_R^q (1, 3, -\frac{2}{3}, \bar{6}_c) \quad (3.8)$$

This model is in some sense left-right symmetric extension of the $SU(2)_L \times U(1)$ model with B-L broken global symmetry, as discussed in the previous section. This, of course, will lead to the existence of $n-\bar{n}$ and $H-\bar{H}$ transitions. The main reason is the existence of an extra term in the Yukawa interaction

$$\Delta L_Y^q = f_e (Q_L^T C \tau_2 \Delta_L^q Q_L + Q_R^T C \tau_2 \Delta_R^q Q_R) + h.c. \quad (3.9)$$

$n-\bar{n}$ oscillators: For M_R which is not astronomically large (which is of phenomenological interest, anyway), the dominant contribution to this process comes from the exchange of the right-handed fields, in the manner depicted in Fig. 2. Assuming

$$m_{\Delta_R^q} = M_R \text{ and using } \langle \Delta_R \rangle = M_R, \text{ one easily obtains}$$

$$T_{n-\bar{n}}^{-1} = \frac{\lambda f_q^3}{M_R^6} \langle \Delta_R \rangle |\psi(0)|^4 = \frac{\lambda f_q^3}{M_R^5} |\psi(0)|^4 \quad (3.10)$$

If λ and f_q are not terribly small, say $\lambda, f_q \sim 10^{-2}$, we get

$$T_{n-\bar{n}} \approx 10^{-13} M_R (\text{GeV})^5 \text{ sec} \quad (3.11)$$

which would require $M_R \approx 10^4$ GeV (or so), in order to make $n-\bar{n}$ transitions observable ($T_{n-\bar{n}} \approx 10^7 - 10^8$ sec).

H-H transition: Here, clearly both Δ_L^q and Δ_R^q exchange leads to H-H transition amplitude. Without fine tuning the parameters, we expect (see Section IV)

$$m_{\Delta_L}^q \sim m_{\Delta_R}^q \sim m_{\Delta_L}^{\ell} \sim m_{\Delta_R}^{\ell} \approx M_R.$$

$$T_{H-\bar{H}}^{-1} \approx \frac{\lambda f_q^3 f_e}{M_R^8} |\psi(0)|^4 \frac{(m_e \alpha)^3}{\pi} \quad (3.12)$$

We have a simple connection between $T_{n-\bar{n}}$ and $T_{H-\bar{H}}$

$$\frac{T_{H-\bar{H}}}{T_{n-\bar{n}}} \approx \frac{\pi M_R^3}{(m_e \alpha)^3 f_e} \quad (3.13)$$

which for $M_R \approx 10^6$ GeV, would make H-H transition hopelessly small: $T_{H-\bar{H}} > 10^{28} T_{n-\bar{n}}$. It is noteworthy that even if M_R is fairly low (~ 100 GeV or so), from (3.13) $T_{H-\bar{H}} > 10^{22} T_{n-\bar{n}}$ and so H-H transition would still be unobservable.

In conclusion, left-right symmetric models based on $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ provide a simple framework in which B-L becomes a dynamic variable, a part of a gauge group. Although the model predicts both $n-\bar{n}$ and H-H, the precise estimates of $T_{n-\bar{n}}$ and $T_{H-\bar{H}}$ are plagued by huge uncertainties. On the other hand, similarly as in the $SU(2)_L \times U(1)$ model with B-L a broken global symmetry, the ratio of $T_{H-\bar{H}}/T_{n-\bar{n}}$ becomes very much independent of all the unknown parameters. However unlike in the previous case, left-right model suggest possibly observable $n-\bar{n}$ oscillations, but not H-H transitions; the latter lifetime being hopelessly large: $T_{H-\bar{H}} > 10^{22}$ yr, from the existing limit $T_{n-\bar{n}} \geq 1$ yr.

IV. GRAND UNIFICATION AND MATTER OSCILLATIONS: A LESSON ON SURVIVAL

As I emphasized in the introduction, minimal SU(5) theory, or any other theory that breaks in the same manner necessarily predicts a desert between M_W and $M_X \approx 10^{14}-10^{15}$ GeV. In particular it makes n - n and H - H transitions, if at all possible, hopelessly unobservable. In the last two sections we have shown how simple it is to find an oasis in this desert, if one is willing to restrict oneself to phenomenological theories such as $SU(2)_L \times U(1)$ or $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$. Matter oscillations, especially in n - n , can easily be made observable and furthermore the oscillation times $T_{n-\bar{n}}$ and $T_{H-\bar{H}}$ become related to each other. The question is, could we expect the same in simple GUTS which are extended beyond SU(5)? The answer, as we shall see, appears to be no, if no fine tuning of the parameters of the theory is allowed. There may or may not be an oasis in the GUT desert, but matter oscillations are unlikely to be observed.

Now, in the simple models these processes are mediated by Higgs scalars, so that the relevant issue becomes whether such Higgs scalars can have naturally reasonably small masses, i.e. below $\sim 10^5$ GeV. I want to give a brief report here on the work done in collaboration with Rabi Mohapatra regarding Higgs mass scales, which completes the earlier work of Georgi²⁸ and del Aquila and Ibanez²⁹. An interested reader can find more detailed discussion in Ref. 27 and in my earlier review³⁰.

To spare the impatient reader unnecessary boredom, let me announce the results before discussing them. Higgs bosons have trouble surviving from becoming superheavy²⁷⁻²⁹, and in particular the exotic scalars which radiate matter oscillations get the mass proportional to M_X . In other words, even when there is an oasis in the desert, they will never know about it. Instead of a general discussion, this will be presented through a simple example²⁷.

Illustrative example: SU(5) with two 5's: The model contains two 5's: H_1 and H_2 and a 24 dimensional multiplet Σ . Imagine the following discrete symmetries implemented for simplicity

$$\begin{array}{ll} D_1: \Sigma \rightarrow -\Sigma & D_2: \Sigma \rightarrow \Sigma \\ H_2 \rightarrow H_1, H_2 \rightarrow H_2 & H_1 \rightarrow H_2 \end{array}$$

The most general Higgs potential is then given by

$$\begin{aligned}
V(\Sigma, H_1, H_2) = & -\frac{1}{2} \mu_X^2 \text{Tr} \Sigma^2 + \frac{1}{4} a (\text{Tr} \Sigma^2)^2 + \frac{1}{2} b \text{Tr} \Sigma^4 \\
& -\frac{1}{2} \mu_W^2 (H_1^\dagger H_1 + H_2^\dagger H_2) + \frac{1}{4} \lambda_1 [(H_1^\dagger H_1)^2 + (H_2^\dagger H_2)^2] + \frac{1}{8} \lambda_3 (H_1^\dagger H_1) (H_2^\dagger H_2) \\
& + \frac{1}{4} \lambda_4 (H_1^\dagger H_2) (H_2^\dagger H_1) + \frac{1}{2} \lambda_5 [(H_1^\dagger H_2)^2 + (H_2^\dagger H_1)^2] + \frac{1}{8} \lambda_6 (H_1^\dagger H_1 + H_2^\dagger H_2) \times (H_1^\dagger H_2 + H_2^\dagger H_1) \\
& + \alpha (H_1^\dagger H_1 + H_2^\dagger H_2) \text{Tr} \Sigma^2 + \beta (H_1^\dagger \Sigma^2 H_1 + H_2^\dagger \Sigma^2 H_2) \\
& + \gamma (H_1^\dagger \Sigma^2 H_2 + H_2^\dagger \Sigma^2 H_1) + \delta (H_1^\dagger H_2 + H_2^\dagger H_1) \text{Tr} \Sigma^2
\end{aligned} \tag{4.2}$$

The diagonalization of the potential leads then to the following values for Higgs boson masses. First, as expected, Σ is superheavy. Second, if we write

$$H_i = \begin{pmatrix} h_i \\ \phi_i \end{pmatrix}, \quad i = 1, 2 \tag{4.3}$$

where h_i are color triplets ($SU(2)_L$ singlets) and ϕ_i are the usual $SU(2)_L$ doublets (color singlets), then h_i and h_2 become superheavy, too, as they should be for the sake of proton stability. This has not taught us much, since there was no freedom up to now. What about the doublets ϕ_i ? Let us define

$$\eta = \sqrt{2} (\phi_1 + \phi_2) \quad ; \quad \rho = \sqrt{2} (\phi_1 - \phi_2) \tag{4.4}$$

which become physical states

$$\eta = \frac{1}{\sqrt{2}} \begin{pmatrix} \sqrt{2} G_W^+ \\ \eta + v_W + i G_Z \end{pmatrix}, \quad \rho = \frac{1}{\sqrt{2}} \begin{pmatrix} \rho^+ \\ \rho_1 + i \rho_2 \end{pmatrix} \tag{4.5}$$

with

$$\langle \phi_1 \rangle = \langle \phi_2 \rangle = \frac{1}{\sqrt{2}} \begin{pmatrix} 0 \\ v_W \end{pmatrix} \quad \text{and } G_W^+ \text{ and } G_Z \text{ are unphysical Higgs bosons}$$

present in the renormalizable gauge which in the unitary gauge get eaten by W^+ and Z , respectively. Clearly then, η must be light ($\sim M_W$) to preserve $SU(2)_L$ invariance. A computation confirms it^{2/}

$$m_n^2 = \left[\frac{1}{4} (\lambda + \lambda_3 + \lambda + \lambda_5 + \lambda_6) - \frac{6}{5} \frac{(10\alpha + 10\delta + 3\beta + 3\gamma)^2}{15a + 7b} - \frac{8}{10} \frac{\beta + \gamma}{6} \right] v_W^2 \quad (4.6)$$

On the other hand, the other doublet ρ has no reason to be light, since $\langle \rho \rangle = 0$. Its mass is a test of a more general issue: what are the masses of the particles which are not phenomenologically forced to be superheavy? Now, no symmetry forbids ρ to receive a superlarge mass and so it does

$$m_\rho^2 = -\frac{15}{2} (d + \frac{3}{10} \gamma) v_X^2 \quad (4.7)$$

We have analyzed²⁷ many other models with the same conclusion, that the survival principle is operative²⁸⁻²⁹: if the minimal fine tuning is assumed, then, unless protected by some discrete, global or gauge symmetry, the Higgs bosons become superheavy.

In particular, those scalars with quantum numbers of di-quarks and di-leptons have no protective symmetry, and so their mass is $\sim M_X$, independently of the existence of intermediate mass scales. To be more precise, at least some of them become superheavy and they suppress matter oscillations.

Before concluding this section, let us, if only briefly discuss the possibilities of observable $n-n$ in transitions in partially or grand unified models. Now that we know that the Higgs bosons tend to pick up the mass of the order M_X , the hope is to be able to find models with low unification scales. This, although not found in simple GUTS based on small groups is still possible in some, more complicated theories³¹⁻³² such as $[SU(6)]^4$ or $SU(16)$. Such models are too complex to be discussed here.

Another alternative would be a partially confined theory, such as the $SU(2)_L \times SU(2)_R \times SU(4)_C$ model of Pati and Salam². Again, the scale of partial unification turns out to be too large: $M_C > 10^{10}$ GeV, from the $\sin^2 \theta_W$ considerations³³. Buras et al., suggest³³ extending the gauge group to $[SU(2)_L \times SU(2)_R]^2 \times SU(4)_C$ in order to bring M_C down; a solution I do not find particularly appealing, although it is certainly legitimate.

A possible hope arises when $SU(2)_L \times SU(2)_R \times SU(4)_C$ is embedded in $SO(10)$. Namely, a possible chain of symmetry breaking³⁴

$$\begin{aligned} O(10) &\xrightarrow{M_X} SU(2)_L \times U(1)_R \times SU(4)_C \xrightarrow{M_C} SU(2)_L \times U(1)_R \times U(1)_{B-L} \times SU(3)_C \\ &\xrightarrow{M_{BL}} SU(2)_L \times U(1) \times SU(3)_C \xrightarrow{M_W} U(1)_{em} \times SU(3)_C \end{aligned} \quad (4.8)$$

allows a solution³⁵ $m_c = 10^5 - 10^6$ GeV, close to its lower bound³⁶ $\sim 10^5$ GeV. However, still not all the relevant lepto-quarks remain at M_c and $n-\bar{n}$, $H-H$ transitions are still suppressed by $1/M_X^2$. Namely, only those with $T_{3R} = 0$ have mass $\sim M_c$, but they are not sufficient to mediate matter oscillations (this can be easily checked by inspection).

V. SUMMARY

In this talk, I have tried to go through the general discussion of matter oscillations in gauge theories. As we learned from the introduction and Section II, this question is intimately tied up to the analysis of B-L operator and its breaking. The special position that B-L holds is best illustrated by the facts that the nonperturbative effects in $SU(2)_L \times U(1)$ fail to break it and also that it becomes an automatic symmetry at the leading $\Delta B \neq 0$ operators in the case of a single mass scale beyond M_W .

On the phenomenological level, it is easy to construct models with observable $n-\bar{n}$ and $H-H$ oscillations. Although by no means the only example, extended $SU(2)_L \times U(1) \times SU(3)_C$ model with B-L spontaneously broken provides an interesting link between $T_{n-\bar{n}}$ and T_{H-H} and neutrino mass; all possibly observable. When the model is made left-right symmetric, B-L becomes gauged; Majoran disappears and $H-H$ oscillations become hopelessly suppressed.

Under the conditions of naturalness, the grand unification of these simple models fails. Namely, they can be easily embedded in $SU(5)$ and $SO(10)$, respectively, however the relevant Higgs scalars which mediate matter oscillations become superheavy. Of course, one may abandon the idea of minimal fine tuning, as has been done in most of the models^{16,22}, in which case one is not surprised to find observable $n-\bar{n}$ and $H-H$ transitions^{16,22,37}.

The main question I have tried to answer here is whether we can hope to observe matter oscillations in the near future, i.e. whether there are single theories that predict $T_{n-\bar{n}}$ and T_{H-H} . Unless counting on luck, my feeling, one which I have tried to portray in this talk is that the answer is no. Both $n-\bar{n}$ and $H-H$ transitions are dramatically dependent on the relevant mass scale, usually Higgs mass scales and if anything, these mass scales are expected in simple GUTS to be very large. Still, on the pragmatic level one cannot yet rule out these processes and as usual we should let the experiment decide.

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