Conf-820470--2

BNL 31423

HIGGS MASS SCALES AND MATTER – ANTIMATTER OSCILLATIONS IN GRAND UNIFIED THEORIES *

BNL--31423

DE82 017638

MASTER

Goran Senjanović

Physics Department

Brookhaven National Laboratory

Upton, N.Y. 11973

DISCLAIMER

This report was prepared as an account of work sponsored by an agency of the United States Government. Neither the United States Government nor any agency thereof, nor any of their employees, makes any warranty, express or implied, or assumes any legal liability or responsibility for the accuracy, completeness, or usefulness of any information, apparatus, product, or process disclosed, or represents that its use would not infringe privately owned rights. Reference herein to any specific commercial product, process, or service by trade name, trademark, manufacturer, or otherwise, does not necessarily constitute or imply its endorsement, recommendation, or favoring by the United States Government or any agency thereof. The views and opinions of authors expressed herein do not necessarily state or reflect those of the United States Government or any agency thereof.

Abstract

A general discussion of mass scales in grand unified theories is presented, with special emphasis on Higgs scalars which mediate neutron-antineutron ($n-\bar{n}$) and hydrogen-antihydrogen ($H-\bar{H}$) oscillations. It is shown that the analogue of survival hypothesis for fermions naturally makes such particles superheavy, thus leading to unobservable lifetimes. If this hypothesis is relaxed, an interesting possibility of potentially observable $n-\bar{n}$ and $H-\bar{H}$ transitions, mutually related arises in the context of SU(5) theory with spontaneously broken B-L symmetry.

* Invited talk at the Workshop on Neutron-Antineutron Oscillations, Harvard Univ., April 30 – May 1, 1982.

The submitted manuscript has been authored under contract DE-AC02-76CH00016 with the U.S. Department of Energy. Accordingly, the U.S. Government retains a nonexclusive, royalty-free license to publish or reproduce the published form of this contribution, or allow others to do so, for U.S. Government purposes.

DISTRIBUTION OF THIS DOCUMENT IS UNLIMITED

The very existence of this Workshop on Neutron Oscillations is the best illustration of the reversal of our attitude towards conservation of baryon (B) and lepton (L) numbers in nature. Previously sacred conservation laws are actully viewed unnecessarily as fundamental principles, and more than that, the apparent excess of matter (M) over antimatter ($\bar{M}$) is now taken as a clue that baryon nonconservation has to take place in some physical processes. On the theoretical side, development of grand unified theories, through the original work [1] of Pati and Salam and Georgi and Glashow has provided a consistent framework in which nonconservation of B and L is predicted to occur.

The burning question we are now facing seems to be not whether, but which $\Delta B \neq 0$ processes can be expected to be observed in near future. A few years ago Weinberg and Wilczek and Zee [2] have developed a model independent method of analyzing this question based on the short distance expansion of the leading effective $\Delta B \neq 0$ operators. Let me describe it here, if only briefly [3]. Since the standard $SU(2)_L xU(1)xSU(3)_c$ theory of electroweak and strong interactions has inbuilt into it, at least in perturbation theory, the conservation of baryon number, the observation of $\Delta B \neq 0$ processes would indicate the existence of new, presumably much heavier particles. Integrating over these new degrees of freedom one can then obtain and categorize effective $\Delta B \neq 0$ operators, invariant under $SU(2)_L xU(1)xSU(3)_c$. In the case of a single new mass scale M_X beyond M_W, the lowest dimension operators would dominate, since the others would be suppressed by some power of M_W/M_X. These operators have dimension d=6 and take the form

$$O = \frac{1}{M_X^2}(q q q \ell) \qquad (1)$$

where q and l stand for quark and lepton fields, respectively. Such operators describe B-L conserving nucleon decays: $p \rightarrow \pi^{o} + e^{+}$ and $n \rightarrow \pi^{-} + e^{+}$. The experimental limit $\tau_p \gtrsim 10^{31}$ years requires then $M_X \gtrsim 10^{14}$ GeV. Such enormous mass scale, as Georgi, Quinn and Weinbegg have shown [4], corresponds to the unification energy of weak, electromagnetic and strong couplings and is predicted in the minimal SU(5) theory. It is not suprising that with great anticipation we now await new experimental results that will probe τ_p up to 10^{32} - 10^{33} years and therefore provide a direct and crucial test of simple grand unification.

If above picture with the "desert" between M_W and M_X is correct the higher dimension operators should be hopelessly unobservable. In particular, neutron-antineutron (n-$\bar{n}$) oscillations, the topic of this workshop, are described by the following type of d=9 operators

$$O_{n-\bar{n}} = \frac{1}{M^5}\left(u^T C^{-1} u\right)\left(d^T C^{-1} d\right)\left(d^T C^{-1} d\right) \qquad (2)$$

where again color indices are ignored. Another possible type of matter-antimatter oscillations: hydrogen-antihydrogen (H-$\bar{H}$) is described by even higher dimension operators (d=12)

$$O_{H-\bar{H}} = \frac{1}{M^8}\left(u^T C^{-1} u\right)\left(u^T C^{-1} u\right)\left(d^T C^{-1} d\right)\left(e^T C^{-1} e\right) \qquad (3)$$

If $M=M_X$, these processes would never be observed. The experimental limits on oscillation lifetimes require $\tau_{n-\bar{n}} \gtrsim 10^7$ sec [5] and $\tau_{H-\bar{H}} \gtrsim 10^{10}$ yr [6] suggest $M \gtrsim 100$-10^5 GeV, depending on the couplings.

The main topic of my talk will be devoted then to the discussion of whether matter oscillations (n-$\bar{n}$ and H-$\bar{H}$) are compatible with grand unification, where we are used to the desert between M_W and $M_X \gtrsim 10^{14}$ GeV, at least in the simple theories. We shall take a long walk through a desert searching for an oasis (i.e. intermediate mass scale) and shall have to watch out for possible mirages.

We can easily enumerate the possibilities that can lead to $M=100-10^5$ GeV:

a. Grand unified theories with low unification scale $M=M_X$, with somehow suppressed proton decay?

b. Existence of low mass scales in so called partially unified theories, such as $SU(2)_L xSU(2)_R xSU(4)_c$ model [1] of Pati and Salam?

c. In many models the mediators of $n-\bar{n}$ oscillations [7] and $H-\bar{H}$ oscillations [8] are expected to be Higgs scalars. Could it be that such Higgs scalars have masses in the range $100-10^5$ GeV, and furthermore, can such masses be naturally expected? This issue, as we shall see, leads to a lesson on "survival", with a rather pessimistic outlook.

In what follows, I shall go through all of the above possibilities, especially case c. Most of my discussion will be reserved for simple models, but I shall try to make, again especially in case c., my assumptions as general as possible.

The rest of this paper is organized as follows. In section II we shall go through the possibilities a.,b. and c. and reach the conclusion that neither one works. Even in the case of Higgs scalars, which is the least understood sector of gauge theories, we shall conclude that the requirement of naturalness makes them superheavy. On the more phenomenological level, if one lets Higgs masses be free, then, of course, both $n-\bar{n}$ and $H-\bar{H}$ oscillations may be observable. In section III an SU(5) model with spontaneously broken B-L symmetry, which relates above processes to the neutrino mass and the masses of doubly charged Higgs scalars, is discussed. Finally, section IV offers a summary and comments.

II. MASS SCALES IN UNIFIED MODELS

a. GUTS with low unification scale.

Simplest grand unified theories, such as SU(5) and SO(10) contain exotic gauge bosons X which couple to both ql and $\overline{qq}$ channels, and therefore directly mediate proton decay. The stability of proton requires $M_X \gtrsim 10^{14}$ GeV. Low unification scale clearly needs more complicated theories. They do exist on the market; the example being the $SU(6)_L \times SU(6)_R \times SU(6)'_L \times SU(6)'_R$ left-right and color-flavor symmetric model [9]. It turns out [9] that in this model $M_X = 10^4$ GeV; however, both nucleon decay and n-$\bar{n}$ oscillations take very complicated forms and I will not discuss them here. I will rather conclude that simple GUTS, presentable to my experimental collegues, are incompatible with low unification scale.

b. Partially unified models.

As a general prototype, let me discuss simple and popular such model based on left-right symmetric $SU(2)_L \times SU(2)_R \times SU(4)_c$ group of Pati and Salam. This model was used by Mohapatra and Marshak [7] in their discussion of n-$\bar{n}$ oscillations. Let us imagine the simplest chain of symmetry breaking

$$SU(2)_L \times SU(2)_R \times SU(4)_c \xrightarrow[M_c]{} SU(2)_L \times U(1) \times SU(3)_c \xrightarrow[M_W]{} U(1)_{em} \times SU(3)_c$$

It would appear that M_c could be low as we wish, since $SU(2)_L$ and $SU(3)_c$ never get unified. However, it is not true [10], since $U(1)_{B-L}$ and $SU(3)_c$ couplings $\bar{\alpha}$ and α_c become equal at M_c: $\bar{\alpha}(M_c) = 3/2\ \alpha_c(M_c)$ (3/2 is just the normalization condition), and similarly $SU(2)_L$ and $SU(2)_R$ couplings: $\alpha_L(M_c) = \alpha_R(M_c)$. From the renormalization group equations for U(1) coupling α', α_L and α_c

$$\frac{1}{\alpha'(M_W)} = \frac{1}{\alpha_R(M_c)} + \frac{1}{\bar{\alpha}(M_c)} - \frac{1}{2\pi} b' \ln \frac{M_c}{M_W}$$

$$\frac{1}{\alpha_L(M_W)} = \frac{1}{\alpha_L(M_c)} - \frac{1}{2\pi} b_L \ln \frac{M_c}{M_W}$$

$$\frac{1}{\alpha_c(M_W)} = \frac{1}{\alpha_c(M_c)} - \frac{1}{2\pi} b_c \ln \frac{M_c}{M_W} \tag{4}$$

with

$$b' = -\frac{20}{9} n_g - \frac{1}{6} N_H; \quad b_L = \frac{22}{3} - \frac{4}{3} n_g - \frac{1}{6} N_H; \quad b_c = 11 - \frac{4}{3} n_g \tag{5}$$

for n_g generations of fermions and N_H light Higgs doublets ($m \sim M_W$), one gets for the M_c scale

$$M_c = M_W \exp\left[\frac{3\pi}{22}\left(1 - 2\sin^2\theta_W(M_W) - \frac{2}{3}\frac{\alpha(M_W)}{\alpha_c(M_W)}\right)\right] \tag{6}$$

Although we cannot calculate $\sin^2\theta_W(M_W)$ as in SU(5) theory, we can still use the experimental limit $\sin^2\theta_W(M_W) \lesssim 0.25$, to set the bound on M_c: $M_c \gtrsim 10^{12}$GeV ($\alpha^{-1}(M_W)$ is the electromagnetic coupling $\simeq 127.5$ and $\alpha_c(M_W) \gtrsim 0.1$).

Similarly as in the case of GUTS, simple partially unified model still leads to the large mass scale above M_W [11], by far too large to corespond to M-$\bar{M}$ transitions. We seem to have reached the point at which a sensible believer in unification gives up the possibility of observing matter oscillations. In desperation, however, we shall try our last hope: maybe the exotic Higgs particles, the possible mediators of n-$\bar{n}$ and H-$\bar{H}$ processes, remain fairly light ($\lesssim 10^5$GeV), despite the enormous mass scales we have been encountering. A purist may be unhappy with the philosophy of attributing new physics to the least understood and rather obscure Higgs sector of gauge theories; nevertheless, we have to pursue it, since it is logically possible. Most of the models on the market that deal with n-$\bar{n}$ oscillations actually just <u>assume</u> that these Higgses are light and then deal with the consequences of such an assumption; here, we shall rather try to spell out the most natural expectation for Higgs mass scales in GUTS.

Before going into an analysis, let me first give a simple example of possible Higgs multiplets that can generate n-$\bar{n}$ and H-$\bar{H}$ transitions. Imagine $SU(2)_L \times U(1) \times SU(3)_c$ gauge theory with the standard fermionic and Higgs assignment <u>and</u> with the additional set of Higgs fields

$$\Delta_\ell\,(3,\,2,\,1_c) \quad ; \quad \Delta_q\,(3,\,-2/3,\,6_c) \qquad (7)$$

where the number in brackets denote $SU(2)_L$, U(1) and $SU(3)_c$ representation content, respectively. This model has been recently discussed in detail by Mohapatra and myself [8,12]. I will come back to it in section III when I discuss its embedding into SU(5). For the moment, let us just notice that the exchange of Higgs scalars from above multiplets with di-lepton and di-quark quantum numbers can lead to $n-\bar{n}$ (Fig.1) and $H-\bar{H}$ process (Fig.2). In the context of our analysis, it then becomes necessary to know the expected masses for Δ_1 and Δ_q, when they are embedded in GUTS.

To spare the impatient reader of unnecessary anticipation, let me first announce the results before discussing them. As it turns out, Higgs boson multiplets have trouble surviving from becoming superheavy; they are able to do so only when they are responsible for symmetry breaking, i.e. when they contain Goldstone bosons. In particular our exotic Higgs bosons which could mediate matter oscillations get the mass proportional to M_X; i.e. independently of the existence of an oasis in the desert, they are expected <u>not</u> to belong to it.

<u>c. Higgs mass scales in GUTS.</u>

The original work in this section, and in the rest of this paper, has been done in colaboration with Rabi Mohapatra [13]. Some of the material in this section was probably known to some other people in the field and some of our assumptions were stated before by del Aguila and Ibanez [14]; however, to the best of my knowledge, the full discussion as given here and in ref.13, has not appeared in literature before.

In order to be able to present a general discussion without alienating an unlearned reader, let me quickly go through a reminder example: minimal SU(5) model of Georgi and Glashow [1].

(i) Higgs mass spectrum in minimal SU(5) model

The Higgs spectrum consists of the 24 dimensional adjoint representation Σ_{24} responsible for superstrong symmetry breaking and the 5 dimensional vector multiplet responsible for breaking of $SU(2)_L xU(1)$ symmetry.

The most general Higgs potential is [15]

$$
\begin{aligned}
V(\Sigma, H) = & -\frac{1}{2}\mu_x^2 \, T_r \Sigma^2 + \frac{1}{4} a (T_r \Sigma^2)^2 + \frac{1}{2} b \, T_r \Sigma^4 \\
& -\frac{1}{2}\mu_\nu^2 \, H^+H + \frac{1}{4}\lambda (H^+H)^2 \\
& + \alpha \, H^+H \, T_r \Sigma^2 + \beta \, H^+ \Sigma^2 H \qquad (8)
\end{aligned}
$$

where we have forbidden the triplic couplings for simplicity ($\Sigma \to -\Sigma$ symmetry). The maximum of the potential is given by [15]

$$
\langle \Sigma \rangle = V_x \begin{pmatrix} 1 & 0 & 0 & 0 & 0 \\ 0 & 1 & 0 & 0 & 0 \\ 0 & 0 & 1 & 0 & 0 \\ 0 & 0 & 0 & -\frac{3}{2} & 0 \\ 0 & 0 & 0 & 0 & -\frac{3}{2} \end{pmatrix}
$$

$$
\langle H \rangle = \frac{V_W}{\sqrt{2}} \begin{pmatrix} 0 \\ 0 \\ 0 \\ 0 \\ 1 \end{pmatrix} \qquad (9)
$$

Since $M_W = gv_W$ and $M_X = gv_X$, we can see the root of the problem: in order to have $v_W/v_X \lesssim 10^{-12}$ as dictated by physics, we have to restore to an incredible fine tuning of the parameters in (8) - this is the infamous hierarchy problem in GUTS. We will offer no new wisdom regarding this problem, but rather pregmatically assume that such fine tuning is done and will be eventually understood - what happens then?. As expected, the members of Σ become superheavy, but what about the H multiplet? Let us write it as

$$H = \begin{pmatrix} h \\ \varphi \end{pmatrix} \tag{10}$$

where h is colored triplet and $SU(2)_L$ singlet and φ is the usual $SU(2)_L xU(1)$ doublet: $\varphi = \begin{pmatrix} \varphi^+ \\ \varphi^0 \end{pmatrix}$. In the renormalizable R gauge, used in performing computations in gauge theories, φ is decomposed into

$$\varphi = \frac{1}{\sqrt{2}} \begin{pmatrix} \sqrt{2}\, G^+ \\ \eta + v_W + i\, G_Z \end{pmatrix} \tag{11}$$

where G^+ with $m_{G^+} = M_W$ and G_Z with $m_{G_Z} = M_Z$ are the Goldstone bosons and η is the physical Higgs scalar. $SU(2)_L$ symmetry then dictates $m_\eta = O(M_W)$; the computation confirms the expectation: $m_\eta = 1/2\, \lambda\, v_W^2$. What about h?. It has no reason to remain light and so it doesn't [15]: $m_h^2 = -5/2\, \beta\, v_X^2$. This is actually a blessing, since h mediates proton decay.

As Georgi [16] has suggested for the fermions, particles do not survive to low energies, if no symmetry forbids their possible large mass term. Naively, it appears that the light doublet violates the survival hypothesis, since it could be given $SU(2)_L xU(1)$ invariant mass $\sim M_X$; however, the Goldstone bosons in the multiplet forbid it.

Now, the reader may and should argue that the example we have given is trivial. The light doublet had to be light, and the colored triplet h had to be heavy for the sake of the stability of the proton, so we have learned nothing about Higgs multiplets whose mass scale is not dictated by physical arguments. A more nontrivial example is furnished by

(ii) SU(5) model with two 5's (H_1 and H_2)

The most general Higgs potential

$$
\begin{aligned}
V(\Sigma, H_1, H_2) = &-\frac{1}{2}\mu_X^2 \, T_r \Sigma^2 + \frac{1}{4} a \left(T_r \Sigma^2\right)^2 + \frac{1}{2} b \, T_r \Sigma^4 \\
&-\frac{1}{2}\mu_W^2 \left(H_1^+ H_1 + H_2^+ H_2\right) + \frac{1}{4}\lambda\left[\left(H_1^+ H_1\right)^2 + \left(H_2^+ H_2\right)^2\right] \\
&+ \frac{1}{2}\lambda_3 \left(H_1^+ H_1\right)\left(H_2^+ H_2\right) + \frac{1}{4}\lambda_4 \left(H_1^+ H_2\right)\left(H_2^+ H_1\right) \\
&+ \frac{1}{2}\lambda_5 \left[\left(H_1^+ H_2\right)^2 + \left(H_2^+ H_1\right)^2\right] + \frac{1}{8}\lambda_6 \left(H_1^+ H_2\right)\left(H_2^+ H_1\right) \\
&+ \alpha \left(H_1^+ H_1 + H_2^+ H_2\right) T_r \Sigma^2 + \beta \left(H_1^+ \Sigma^2 H_1 + H_2^+ \Sigma^2 H_2\right) \\
&+ \gamma \left(H_1^+ \Sigma^2 H_2 + H_2^+ \Sigma^2 H_1\right) + \delta \left(H_1^+ H_2 + H_2^+ H_1\right) T_r \Sigma^2
\end{aligned}
\tag{12}
$$

is invariant under $SU(5)\times D_1 \times D_2$, with

$$
\begin{array}{ll}
D_1: \ \Sigma \to -\Sigma & D_2: \ \Sigma \to \Sigma \\
\quad H_a \to H_a \ (a=1,2) & \quad H_1 \leftrightarrow H_2
\end{array}
\tag{13}
$$

The discrete symmetries were assumed for the sake of simplicity, leading to no loss of generality. The diagonalization of the potential leads then to the following v alues for Higgs boson masses. First, as before, Σ is superheavy. Second, h_1 and h_2 become superheavy, as they should for the sake of proton stability. This is not an useful information; we have no freedom in this case. What about the doublets φ_1 and φ_2?. Let us define the physical states

$$
\eta = \frac{1}{\sqrt{2}} \left(\varphi_1 + \varphi_2 \right) \ ; \quad \rho = \frac{1}{\sqrt{2}} \left(\varphi_1 - \varphi_2 \right)
\tag{14}
$$

It is easy to see that [13]

$$
\eta = \frac{1}{\sqrt{2}} \begin{pmatrix} \sqrt{2}\, G_W^+ \\ \eta + V_W + i G_Z \end{pmatrix}, \quad \rho = \begin{pmatrix} \rho^+ \\ \rho_1 + i \rho_2 \end{pmatrix}
\tag{15}
$$

where $\langle \varphi \rangle_1 = \langle \varphi \rangle_2 = 1/\ 2 \begin{pmatrix} 0 \\ v_W \end{pmatrix}$ is the CP conserving minimum of the potential.

Since η contains Goldstone bosons it will be light. A computation gives [13]

$$
m_\eta^2 = \left[\frac{1}{4} \left(\lambda + \lambda_3 + \lambda_4 + \lambda_5 + \lambda_6 \right) - \frac{6}{5} \frac{\left(10\alpha + 10\delta + 3\beta + 3\gamma\right)^2}{15a + 7b} - \frac{9}{10} \frac{\beta + \gamma}{b} \right] V_W^2
\tag{16}
$$

The other doublet ρ has no Goldstone bosons and so no symmetry forbids it to receive a superlarge mass. Again, physical intuition does not fail; we find $SU(2)\times U(1)$ invariant mass term

$$m_\Sigma^2 = -\frac{15}{2}\left(\delta + \frac{3}{10}\gamma\right) V_X^2 + O(V_W^2) \qquad (17)$$

We have analyzed many other models [13], and always ended with some conclusion as above. We can suggest a precise definition of the extension of survival hypothesis.

(iii) General rules for Higgs boson masses.

1. When a multiplet of Higgs scalars Φ_S, being a representation of a subgroup G_S of the original gauge group G, is used to break the symmetry through its vacuum expectation value, i.e. when $G_S \xrightarrow[\langle\Phi_S\rangle]{} G'_S$, then all the members of Φ_S multiplet have the mass $\sim \langle\Phi_S\rangle$.

The example of this statement is the "physical" doublet of $SU(2)_L xU(1)$ subgroup of SU(5); since $SU(2)_L xU(1) \xrightarrow[\langle\eta\rangle]{} U(1)_{em}$, we get $m_\eta \sim \langle\eta\rangle = M_W$. The reason for this is the existence of Goldstone bosons, as we have stressed before.

2. If minimal fine tuning is assumed (needed for 1.), then, unless protected by some discrete or global (either accidental or not), all other Higgs particles will become superheavy.

In particular, Higgs scalars with quantum numbers of diquarks and di-leptons have no symmetry to prevent them from becoming superheavy; they join the rest of the victims who fail to survive to low energies. To be precise, the particles of our interest will have the mass $\sim (\sum\lambda_i)v_X$, where $\sum\lambda_i$ is some combination of scalar self-couplings. This is true, even when there are intermediate mass scales, corresponding to new electroweak thresholds such as say $SU(2)_L xSU(2)_R xU(1)$ in O(10) model; still the relevant Higgs bosons get the mass proportional to M_X. Matter oscillations, although allowed in simple grand unified theories, are expected (if the principle of naturalness hold through) to be dramatically suppressed: $\tau_{M-\bar{M}} > 10^{30}$ years, which would take them out of reach of experiments.

Does this mean that we have rigorously proved unobservability of matter oscillations in GUTS?. Unfortunately (or fortunately?) the answer is no. Recall that in order to make GUTS consistent we have to restore to a dirty procedure of incredible fine tuning of parameters ($v_W/v_X \lesssim 10^{-12}$). The devil's advocate may argue that the same procedure may be repeated in the case of exotic Higgses: although $m_H = (\sum \lambda_i)v_X$, we can make $\sum \lambda_i$ arbitrarily small to ensure such Higgses to be as light as we wish. Although not supported by physical arguments, from the mathematical point of view this is no worse than what we did for the original gauge hierarchy.

In the next section I will describe a GUT model of n-$\bar{n}$ and H-$\bar{H}$ oscillations, in which one gives the only minimal fine tuning requirement. One still finds some restrictions on the model, despite the increased freedom of Higgs masses.

III. SPONTANEOUS BREAKDOWN OF GLOBAL B-L SYMMETRY AND MATTER OSCILLATIONS IN SU(5)

In section II we have briefly mentioned a recently suggested model [8,12] of n-$\bar{n}$ and H-$\bar{H}$ transitions. I will discuss here its grand unified version [17], with the special emphasis paid to the predictions for proton lifetime and $\sin^2\theta_W$, besides matter oscillations.

Gauge group:

$$G = SU(5)xU(1)_G$$

with $U(1)_G$ a global symmetry specified below.

Fermions:

As in the minimal SU(5) model, we have

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

$$\overline{\underline{10}} = T_L = \begin{pmatrix} 0 & u_3^c & -u_2^c & -u_1 & -d_1 \\ -u_3^c & 0 & u_1^c & -u_2 & -d_2 \\ u_2^c & -u_1^c & 0 & -u_3 & -d_3 \\ u_1 & u_2 & u_3 & 0 & e^+ \\ d_1 & d_2 & d_3 & -e^+ & 0 \end{pmatrix}_L \qquad (18)$$

Higgs:

We have Σ_{24} and H_5 as in the minimal model, with the addition of $\underline{15}$ dimensional multiplet Δ_{15} and $\underline{50}$ dimensional representation χ_{50}. The last multiplet was used recently by Arnellos and Marciano [18] in order to generate H-$\bar{H}$ transitions.

The most general Yukawa interaction consistent with SU(5)x $U(1)_G$ symmetry is

$$L_Y = h_1 \bar{F}_R H^* T_L + h_2 T_L^T C^{-1} H T_L$$

$$+ h_\Delta F_R^T C^{-1} \Delta^* F_R + h_\chi T_L^T C^{-1} \chi T_L + h.c. \qquad (19)$$

where we have suppressed SU(5) indices. The transformation properties of the fields under $U(1)_G$ are

$\underline{U(1)_G}$:

$$T_L \to e^{i\theta} T_L \; ; \quad F_R \to e^{3i\theta} F_R \; ; \quad \Sigma \to \Sigma$$

$$H \to e^{-2i\theta} H, \quad \Delta \to e^{6i\theta} \Delta, \quad \chi \to e^{-2i\theta} \chi \qquad (20)$$

Succession $\langle \Sigma \rangle \neq 0$, $\langle H \rangle \neq 0$ breaks $SU(5) \to U(1)_{em} \times SU(3)_c$. What about $U(1)_G$ symmetry? $\langle \Sigma \rangle \neq 0$ cannot break; but $G \langle H \rangle \neq 0$; however $(Y+G/2)\langle H \rangle = 0$ (G is the generator of $U(1)_G$). A simple computation gives [19]

$$B - L = \frac{2}{5}\left(Y + \frac{G}{2} \right) \qquad (21)$$

so that B-L is still conserved global symmetry. It will be spontaneously broken through $\langle \Delta \rangle \neq 0$. The main manifestation of B-L breaking is a Majorana neutrino mass [20]

$$m_\nu = h_\Delta \langle \Delta \rangle \qquad (22)$$

Another crucial feature of the theory is the existence of a Goldstone boson, a Majoron [21], which is mainly coupled to neutrinos. Its properties have been discussed at length before [22]; for us it will only be important to keep in mind a constraint $h \leq 7 \times 10^{-3}$.

<u>n-$\bar{n}$ oscillations</u>

The amount of neutron oscillations, being a $\Delta B = 2$, $\Delta L = 0$

process, is necessarily proportional to a measure of B-L symmetry breaking, i.e. neutrino mass. The $SU(5)xU(1)_G$ symmetry allows the $\underline{50}\ \underline{50}\ \underline{50}\ \underline{15}$ coupling in the potential $V(\chi, \Delta) = \lambda \chi^3 \Delta$, leading to a graph of the type in Fig.1, where three χ fields meet vacuum insertion $\langle \Delta \rangle$. It is a simple exercise to estimate the lifetime $\tau_{n-\bar{n}}$ for neutron oscillation

$$\tau^{-1}_{n-\bar{n}} = \lambda \frac{h_\chi^3 \, m_\nu}{m_\chi^6 \, h_\Delta} |\psi(0)|^4 \tag{23}$$

where we have used (22); $|\psi(0)|^4$ is the quark wave function inside a nucleon [23]: $|\psi(0)|^4 \simeq 10^{-3} (\mathrm{GeV})^6$ and we have assumed common Higgs mass m_χ for the colored sextets in χ.

H-$\bar{H}$ oscillations and double proton decay

$H \rightarrow \bar{H}$ ($pe \rightarrow \bar{p}e^+$) transition or double protob decay $pp \rightarrow e^+e^+$ are $\Delta B = 2$, $\Delta(B-L) = 0$ processes and so they proceed in the absence of B-L breaking. There are two different contributions. The first is through the $(\underline{50})^4$ couplings [18] $\lambda_\chi(\chi^+\chi)^2$, which generate the amplitude shown in Fig.2. Due to presence of Δ_{15}, we an additional graph through $\underline{50}\ \underline{50}\ \underline{50}\ \underline{15}$ coupling $\lambda \chi^3 \Delta$. The inverse lifetime for H-$\bar{H}$ transition becomes [6]

$$\tau^{-1}_{H-\bar{H}} \simeq \frac{h_\chi^3}{m_\chi^6} \left(\frac{\lambda_\chi h_\chi}{m^2_{\chi^{++}}} + \frac{\lambda h_\Delta}{m^2_{\Delta^{++}}} \right) |\psi(0)|^4 \frac{(m_e \alpha)^3}{\pi} \tag{24}$$

where χ^{++} and Δ^{++} are doubly charged Higgs fields from χ and Δ, respectively. $\Delta(B-L) = 0$ H-$\bar{H}$ transition is not suppressed by m_ν , but being an operator of higher dimension it is suppressed by $(m_e/m^{++})^2$.

The main problem in predicting n-$\bar{n}$ and H-$\bar{H}$ and $2p \to 2e^+$ lifetimes stems from the enormous uncertainty due to large powers of unknown mass m_χ . However, most of the uncertainty dissapears when the ratio of $\tau_{H-\bar{H}}$ and $\tau_{n-\bar{n}}$ is formed

$$\frac{\tau_{n-\bar{n}}}{\tau_{H-\bar{H}}} \simeq h_\Delta \left(\frac{h_\Delta}{m^2_{\Delta^{++}}} + \frac{\lambda_\chi \, h_\chi}{\lambda \, m^2_{\chi^{++}}} \right) \frac{(m_e \alpha)^3}{\pi m_\nu} \tag{25}$$

We have succeeded, if not in predicting $\tau_{n-\bar{n}}$ and $\tau_{H-\bar{H}}$, at least in relating their ratios to m_ν and the masses of doubly charged Higgs particles. For example, for natural values $\lambda_\chi \simeq \lambda$, $h_\Delta \simeq h_\chi = h$, $m_{\Delta^{++}} \simeq m_{\chi^{++}} = m_{++}$ we would get

$$\frac{\tau_{n-\bar{n}}}{\tau_{H-\bar{H}}} \simeq \frac{h^2 (m_e \alpha)^3}{\pi \, m_\nu \, m^2_{++}} \tag{26}$$

For reasonable values $m_{++} \simeq 20$ GeV, $m_\nu = 10^{-1}$ eV and $h \lesssim 10^{-3}$, the above ratio becomes $\tau_{H-\bar{H}} \simeq 10^{13} \, \tau_{n-\bar{n}}$. Furthermore, light color sextets from χ with $m_\chi \simeq 100$ GeV would give $\tau_{n-\bar{n}} \simeq 10^7$ sec, which would be observable in the experiments now planned[24]. H-$\bar{H}$ oscillation time may be a bit too long ($\gtrsim 10^{13}$ yr) to be observed; however double proton decay with $\tau_{2p} \gtrsim 10^{31}$ yr has a promise of being competitive with the usual $\Delta B = 1$ proton decay.

To make our presentation complete, we have to calculate the effects of additional Higgs multiplets in $\underline{15}$ and $\underline{50}$ on τ_p and $\sin^2\theta_W$. Let us first decompose $\underline{15}$ and $\underline{50}$ into $SU(2)_L \times U(1) \times SU(3)_c$ components

$$\Delta_{15} = \underbrace{(3, 2, 1_c)}_{\Delta_1} + (1, \underbrace{-4/3, 6_c)}_{\Delta_2} + \underbrace{(2, 1/3, 3_c)}_{\Delta_3}$$

$$\chi_{50} = \underbrace{(1, -4, 1_c)}_{\chi_1} + \underbrace{(1, -2/3, 3_c)}_{\chi_2} + \underbrace{(2, -7/3, 3_c)}_{\chi_3}$$

$$+ \underbrace{(3, -2/3, 6_c)}_{\chi_4} + \underbrace{(1, 8/3, \bar{6}_c)}_{\chi_5} + \underbrace{(2, 1, \bar{8}_c)}_{\chi_6} \qquad (27)$$

The approximate change of τ_p due to these new fields is [25]

$$\frac{\bar{\tau}_p}{\tau_p} = \left(\frac{m_{\Delta_1}^5}{m_{\Delta_2}^4 \; m_{\Delta_3}} \cdot \frac{m_{\chi_1}^4 \; m_{\chi_3}^7 \; m_{\chi_5}^4}{m_{\chi_2} \; m_{\chi_4}^6 \; m_{\chi_6}^8} \right)^{4/67} \qquad (28)$$

where $\bar{\tau}_p$ is the new proton lifetime and τ_p is proton lifetime in the minimal SU(5) model. Similarly, one finds for the relative change of $\sin^2\theta_W(M_W)$ compared to the minimal model [25]

$$\Delta \sin^2\theta_W(M_W) = - \frac{\alpha(M_W)}{804\pi} \left[\ln \frac{m_{\Delta_1}^{153} \; m_{\Delta_3}^{90}}{m_{\Delta_2}^{243}} \right.$$

$$\left. + \ln \frac{m_{\chi_4}^{741}}{m_{\chi_1}^{92} \; m_{\chi_2}^{44} \; m_{\chi_3}^{94} \; m_{\chi_5}^{427} \; m_{\chi_6}^{84}} \right] \qquad (29)$$

Except for the case of χ_2 which has to be superheavy since it mediates proton decay, all other masses are unfortunately free parameters, once we give up a principle of minimal fine tuning. It is maybe then natural to reverse the philosophy: maybe for some reason the particles that do not have to be superheavy in order to guarantee the proton stability, somehow remain light. In this case, $m_{\chi_2} = M_X$ and all other masses $\sim M_W$. This gives

$$\bar{\tau}_p = \frac{1}{5}\tau_p \; ; \quad \Delta \sin^2\theta_W \simeq 4\times 10^{-3} \qquad (30)$$

Both effects are tiny and appear to be below the uncertainties in the computation of $\sin^2\theta_W$ and τ_p; the model is in this sense equivalent to minimal SU(5). To be honest, it is possible to change (30) completely, as is obvious from (2?) and (29), if some masses are artificially assigned different values. For example, increasing, say, m_{χ_3} sufficiently could make $\bar{\tau}_p$ large enough to make $\Delta B = 1$ proton decay unobservable; so that we could have an amusing situation of $\Delta B = 2$ $2p \to 2e^+$ decay dominant and possibly observable.

IV. SUMMARY

Simple grand unified theories, such as SU(5) and SO(10), require the existence of an enormous desert between $M_W \simeq 80$ GeV and $M_X \gtrsim 10^{14}$ GeV. We have made an attempt in this talk to see whether the idea of simple grand unification can be made compatible with the possible observability of neutron and hydrogen oscillations. Actually, we have relaxed the restrictions of the theoretical framework by discussing partially unified theories, such as $SU(2)_L \times SU(2)_R \times SU(4)_c$. Still, the desert emerged. The only possibility thus remained that somehow exotic Higgs particles which can mediate the above mentioned processes remain light, despite their natural tendency to receive superlarge masses of order M_X. This, rather phenomenological, approach, although legitimate, requires further, fantastically precise tuning of the parameters. I still feel that this should continue to make us feel rather uneasy about matter oscillations in simple GUTS.

It is often said that living in the real world requires a somewhat pragmatic attitude. That was the essence of our philosophy, after concluding in section II that it tends to contradict the assumption of naturalness in physics, in accepting the possibility of light Higgs scalars which mediate matter oscillations. In section III we have discussed SU(5) model with the spontaneous breakdown of B-L symmetry, that relates potentially observable n-$\bar{n}$ and H-$\bar{\text{H}}$ transitions to neutrino and doubly charged Higgs masses. The model incorporates a possibly fast double proton decay $2p \to 2e^+$ with $\tau_{2p} \gtrsim 10^{31}$ years; and, although it fails in predicting τ_p and $\sin^2\theta_W$, in its simplest and most natural case suggests these values to be close to the minimal SU(5) model predictions.

Acknowledgments

The work described in this talk has been done in collaboration with Rabi Mohapatra. I would also like to thank Maurice Goldhaber, Bill Marciano and Saško Šokorac for discussions.

REFERENCES AND FOOTNOTES

1. J. C. Pati and A. Salam, Phys. Rev. D8, 1240 (1973); ibid. 10, 275 (1974); H. Georgi and S. L. Glashow, Phys. Rev. Lett. 32, 438 (1974).
2. S. Weinberg, Phys. Rev. Lett. 43, 1566 (1979); F. Wilczek and A. Zee, ibid. 43, 1571 (1979).
3. For a nice expose, see A. Zee's talk at this workshop.
4. H. Georgi, H. Quinn and S. Weinberg, Phys. Rev. Lett. 33, 451 (1974).
5. M. Chetyrkin, V. Kuzmin, M. Shaposhnikov and M. Kazanowski, Phys. Lett. 99B, 358 (1981); Riazuddin, Phys. Rev. D25, 885 (1982): D. Peaslee, DOE preprint (1981).
6. G. Feinberg, M. Goldhaber and G. Steigman, Phys. Rev. D18, 1602 (1978). This paper discusses in detail the physics of H-$\bar{H}$ oscillations and double proton decay $p+p \rightarrow e^{+}+e^{+}$.
7. S. L. Glashow, Cargese lectures (1979); R. N. Mohapatra and R. E. Marshak, Phys. Rev. Lett. 44, 1316 (1980); L. N. Chang and N.P. Chang, Phys. Lett. 92B, 103 (1980). For earlier work see V. Kuzmin, Pisma JETP 13, 335 (1970). See also T. K. Kuo and S. T. Love, Phys. Rev. Lett. 45, 93 (1980).
8. R. N. Mohapatra and G . Senjanović, BNL preprint (1982); L. Arnellos and W. J. Marciano, BNL preprint (1982).
9. V. Elias, J. C. Pati and A. Salam, Phys. Rev. Lett. 40, 920 (1978); V. Elias and S. Rajpoot, Phys. Rev. D20, 2445 (1979).
10. V. Elias, Phys. Rev. D14, 1896 (1976); J. D. Bjorken, A. Buras and P. Q. Hung, Fermilab preprint (1981).
11. When $SU(2)_L \times SU(2)_R \times SU(4)_c$ is embedded into SO(10), M_c can be lowered, if $SU(2)_R$ is broken at M_X (as I learned first from

S. Rajpoot and J. Pati). However, as we show in section c., the relevant particles then still get the superlarge mass M_X.

12. See also Mohapatra's talk at this workshop, CCNY preprint (1982).
13. R. N. Mohapatra and G. Senjanović, BNL preprint (to appear).
14. F. del Aguila and L. Ibanez, Nucl. Phys. B177, 60 (1981).
15. A. Buras, J. Ellis, M. K. Gaillard and D. V. Nanopoulos, Nucl. Phys. B135, 66 (1978).
16. H. Georgi, Nucl. Phys. B156, 126 (1979).
17. R. N. Mohapatra and G. Senjanović, BNL preprint (1982).
18. L. Arnellos and W. J. Marciano, ref.8.
19. F. Wilczek and A. Zee, Phys. Lett. (1979).
20. A 15 of Higgs was first used to generate Majorana neutrino mass and $n-\bar{n}$ oscillations by S. L. Glashow, ref.7. However, if only 5 and 15 are present, $n-\bar{n}$ transition rate is hopelessly small due to the large mass of the colored triplet member of 5.
21. Y. Chikashige, R. N. Mohapatra and R. D. Peccei, Phys. Lett. 98B, 265 (1981); G. Gelmini and M. Roncadelli, Phys. Lett. 99B, 411 (1981).
22. H. Georgi, S. L. Glashow and S. Nussinov, Nucl. Phys. B (to appear); V. Barger, W. Y. Keung and S. Pakvasa, Phys. Rev. D25, 907 (1982).
23. For an estimate see, for example, J. Pasupathy, Bangalore preprint (1982).
24. See the proceedings of this workshop.
25. The effects of χ_{50} have been calculated already by Arnellos and Marciano, ref.8.

FIGURE CAPTIONS

Fig. 1 The basic diagram for $n-\bar{n}$ oscillations. The dashed lines represent Higgs bosons which mediate the process and the blob stands for a vacuum expectation value of some Higgs field.

Fig. 2 A typical graph representing $H-\bar{H}$ oscillations and the double proton decay $p+p \longrightarrow e^{+}+e^{+}$.

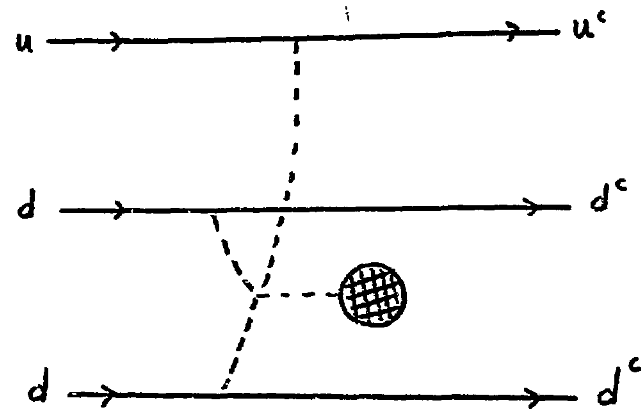

FIG. 1

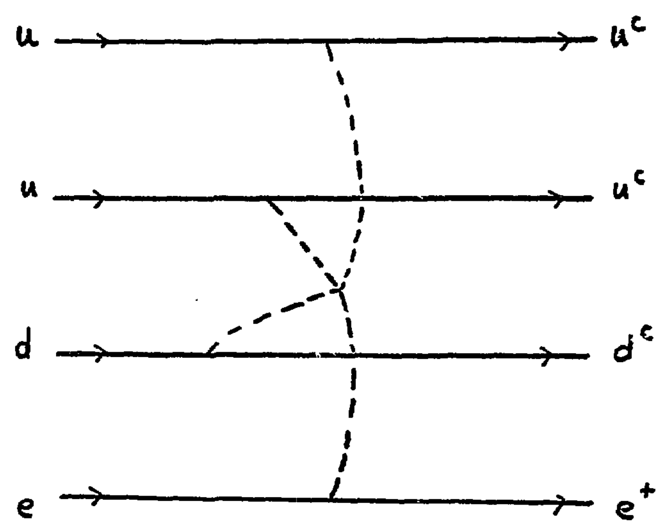

FIG. 2