Neutrino Mass and Baryon Number Non-Conservation in Superstring Models

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Abstract

We propose new mechanisms for understanding neutrino masses in superstring models that contain E_6-singlet zero mass fields after compactification. We show that the low energy gauge group of these models can be phenomenologically acceptable. We then comment on $\Delta B = 1$ and $\Delta B = 2$ baryon number violating processes in these models.

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1. Recently, there has been a great deal of activity in superstring theories with the gauge group $E_8 \times E_8'$.[1,2,3,4,5] The zero slope limit of these theories leads to an anomaly free ten dimensional $E_8 \times E_8'$ Yang-Mills theory coupled to $N = 1$ supergravity. When six extradimensions are compactified[3] to a Calabi-Yau manifold with SU(3)-holonomy, an $N = 1$ locally supersymmetric four dimensional grandunified theory based on gauge group $E_8' \times E_6$ emerges with N_g copies of massless {27} dimensional (under E_6) and $b_{1,1}$ pairs of $\{27\} + \{\overline{27}\}$ chiral superfields (where $b_{1,1}$ is the Bette-Hodge number). By an appropriate choice of[3,4,5] the Calabi-Yau space, one can have $N_g = 3$ or 4. One can then assume the observed matter fields (quarks and leptons) to belong to the {27} -dimensional representations of E_6. This model, therefore, has all the right ingredients for being a candidate theory that unifies all matter and all interactions in nature.[6]

Even though this program of unification appears very attractive, several potential difficulties appear as soon as one starts to study its detailed phenomenological implications: too fast proton decay, potentially large neutrino masses, problem of light Higgs multiplets, and lack of a proper mechanism for supersymmetry breaking. In this paper, we will concern ourselves with only the first three problems. We will exhibit mechanisms for understanding small neutrino masses using light E_6-singlet fields. These models require the existence of light Higgs fields with specific quantum numbers. We show that this requirement can be satisfied for phenomenologically acceptable low energy gauge groups. In some of these models, the $SU(2)_L$ doublet Higgs responsible for symmetry breaking must arise from matter superfields. This leads us to discuss the question of baryon non-conservation such as proton-decay and neutron-anti-neutron oscillation in these models. We comment on the possibility that neutron-anti-neutron oscillation may be observable in this class of models under certain circumstances while avoiding catastrophic proton decay.

To exhibit the problem with the neutrino masses, we note that these theories contain both the left and the right-handed neutrinos; as a result, a large Dirac mass (of order of the charged lepton mass) for neutrinos arises naturally after the breaking of $SU(2)_L \times U(1)_Y$ gauge symmetry. The usual solution for this problem is to employ the "see saw" mechanism,[7] where, by appropriate choice of the Higgs multiplets and symmetry breaking, a large Majorana mass is generated for the right-handed neutrinos. Unfortunately, in superstring models, the Higgs multiplet that can lead to a large Majorana mass for the right-handed neutrino is absent. As a result, new mechanisms have recently been considered to solve this problem.[8,9,10] They can be divided into three categeories: (i) Low scale B-L breaking[8] via non-zero v.e.v. for the superpartner of ν_R ; (ii) Intermediate scale B-L breaking and use of higher dimensional operators;[9] (iii) Use of E_6-singlet Higgs superfields.[10] It is the purpose of this paper to exhibit a detailed model that uses the scenario (iii) and then show that all these models can lead to observable neutron-anti-neutron oscillation, while at the same time avoiding a rapid proton decay.

2. The first difficulty in obtaining an acceptable pattern of neutrino masses in superstring models arises from the presence of new exotic neutral fermions beyond the usual left- and right-handed neutrinos. This can be seen from the decomposition of {27} - dimensional representation of E_6 under the [SO(10), SU(5)] subgroups:

$$
\begin{aligned}
\{27\} = \quad & [16,10] \quad + \quad [16,\bar{5}] \quad + \quad [16,1] \\
& (u,d;u^c,e^c) \quad + \quad (d^c;\nu,e) \quad + \quad \nu^c \\
& + [10,\bar{5}] \quad + \quad [10,5] \quad + \quad [1,1] \\
& (D^c,N,E^-) \quad + \quad [D,N^c,E^+] \quad + \quad n_0
\end{aligned} \tag{1}
$$

The various particles are identified below each group representation. In what follows, we represent a matter multiplet by ϕ and a Higgs multiplet by H. The five neutral leptons are: $(\nu\ ,\ \nu^c,\ N\ ,\ N^c,\ n_o)$. The new neutral leptons $(N,\ N^c)$ must be massive enough so that their contribution to the present energy density of the universe is below the critical density. As far as n_0 is concerned, it can be massless or superheavy depending on whether it couples to superheavy or light gauge bosons. Assuming that N, N^c and n_0 decouple from low energies, we are still left with a Dirac neutrino obtained by combining ν_L and ν_L^c with a mass $m_D \simeq O(m_e) \approx O(\mathrm{MeV})$. The simplest way to avoid this problem would be to set $m_D = 0$ by choosing the appropriate Yukawa coupling to vanish. This would, of course, be, theoretically, unnatural but even phenomenologically, it will put a lower limit the scale of the B-L breaking (or the mass of the second Z'-boson) from considerations of nucleosynthesis,[11] which may be too constraining from the point of view of a "natural" gauge hierarchy.

To solve this problem, we seek ways by which ν^c acquires a large mass. A Majorana mass for ν^c would require breaking (B-L) by two units and is therefore not possible to have at the tree level. They can, however, be induced in either of the following ways:

(i) A higher dimensional term[9] of the form $\frac{1}{M}\{27\}_\phi\ \{\overline{27}\}_H \cdot \{27\}_\phi \cdot \{\overline{27}\}_H$ leads to an effective Majorana mass for the right-handed neutrino ν^c :

$$M_{\nu^c} = V_{BL}^2/M \tag{2}$$

where

$$\langle \tilde{\nu}_H^c \rangle = V_{BL} \ .$$

This leads to a light neutrino mass

$$m_\nu \approx m_D^2 M/V_{BL}^2 \ . \tag{3}$$

Choosing $M \simeq M_{Planck} \simeq 10^{18}$ GeV , $m_D \simeq 1$ MeV requires that $V_{BL} \simeq 10^{11}$ GeV leads to $m_\nu \simeq .1$ eV .

(ii) The second possibility is to use another neutral fermion which is B-L neutral to form a $\Delta(B-L) = 1$ Dirac mass term. This mechanism was used in ref. 8, where the extra neutral fermion chosen was the gaugino corresponding to B-L symmetry. in this paper, we replace the gaugino by an E_6-neutral fermion (denoted by S), that may be present in superstring models. The scenario outlined below realizes this mechanism.

3. In this section, we present two models for neutrino masses. We first outline the details of the models that are relevant only to the discussion of neutrino masses and in a subsequent section discuss the associated low energy gauge group. Both the models we present will require $b_{1,1} = 2$ i.e. two pairs of $\{27\} + \{\overline{27}\}$ representations that act as Higgs fields denoted by H and J respectively. We will then assume that the SO(10) singlet components (denoted by n_0) and [SO(10), SU(5)] representation [16,1] (denoted by ν^c) remain light. The reason for this is that we would like to give them intermediate scale vev (without breaking supersymmetry i.e. maintaining a D-flat direction).

a) Model I:

To write down the most general low energy superpotential in this theory, we first denote the various superfields in the model as follows: using [SO(10), SU(5)] notation: $\left(Q = (u,d)\ ;\ \ell = (\nu, e^-)\right)$

$[16, 10]:\ Q\ ,\ u^c\ ,\ e^c$

$[16, \bar{5}]:\ d^c\ ,\ \ell$

$[16, 1]:\ \nu^c$

$[10, 5]:\ D\ ,\ (N^c, E^+) \equiv E$

$[10, \bar{5}]:\ D^c,\ (N, E^-) = E^c$

$[1, 1]:\ n_o$.

Denoting the components of the Higgs field by a subscript J and H and suppressing all generation indices, we can write the super potential as follows: $P_I = P_o + P_I'$ where

$$\begin{aligned} P_o = {} & \lambda_1 QQD + \lambda_2 Qu^cE + \lambda_3 Qd^cE^c \\ & + \lambda_4 QD^c\ell + \lambda_5 u^c d^c D^c + \lambda_6 u^c D e^c \\ & + \sum_{a=\text{matter},H} \lambda_7^a d^c D \nu_a^c + \sum_{a=\text{matter},H} \lambda_8^a D^c D n_{0,a} + \sum_{a=\text{matter},H} \lambda_9^a \ell E \nu_a^c \\ & + \sum_{a=\text{matter},H} \lambda_{10}^a E E^c n_{0,a} + \lambda_{11} \ell E^c e^c \\ & + \beta_1 \nu_H^c \bar{\nu}_J^c S + \beta_2 n_{0,H} \bar{n}_{0,J} S \end{aligned} \tag{4a}$$

and

$$P'_I = \lambda_{12}\, n_0 \bar{n}_{0,J}\, S + \lambda_{13} \nu^c \bar{\nu}^c_J S \tag{4b}$$

In P'_I , Higgs field H contributions are eliminated by imposing a discrete symmetry under which the E_6-singlet S and $J,\bar{J}$ are odd and all other fields are even. We first assume that the SO(10) singlet component of H acquires a v.e.v. V_6 along a D-flat direction; subsequently, two components of $\bar{J}$, one along the SO(10) singlet and another along the [16,1] (or ν^c direction) acquire the following effective v.e.v.:

$$\langle \bar{J}(n_0) \rangle = \mu \quad \text{and} \quad \langle \bar{J}(\nu^c) \rangle = V_{BL} \tag{5}$$

with

$$\mu \ll V_{BL} \ll V_6 \; .$$

The neutral fermions N and N^c pick up large mass V_6 and decouple from low energies. Also, it gives large mass to the D-quark, which decouples from the low energy sector. We further assume $\beta_1 \ll \lambda_{13}$. Defining, $\nu^{c'} = \nu^c + (\beta_1/\lambda_{13})\, \nu^c_H$ and $n'_o = (\beta_2 n_{OH} + \lambda_{12} n_o)/(\beta_2^2 + \lambda_{12}^2)^{1/2}$, we find (the combinations orthagonal to $\nu^{c'}$ and n'_o remain massless and invisible), the 4×4 mass matrix for the remaining neutral fermions ν , $\nu^{c'}$, n'_o , S , to be of the form:

$$\begin{array}{c|cccc} & \nu & n'_o & \nu^{c'} & S \\ \hline \nu & 0 & 0 & m_D & 0 \\ n'_o & 0 & 0 & 0 & \mu \\ \nu^{c'} & m_D & 0 & 0 & V_{BL} \\ S & 0 & \mu & V_{BL} & 0 \end{array} \; . \tag{6}$$

It is easy to see that on diagonalizing this matrix, we obtain two Dirac particles with masses:

$$m_1 = \frac{m_D \mu}{V_{BL}}$$

and

$$M_2 = V_{BL} \ . \qquad (6a)$$

If we choose $V_{BL} \simeq 10^{12}$ GeV , this implies that for $\mu < 10^6$ GeV , $m_1 < 1$ eV . Thus, in this picture neutrino is a Dirac particle with a naturally small mass. We point out that this is a completely new mechanism for generating light Dirac neutrinos and could be useful in general supsymmetric models. In this model, neutrinoless double beta decay will be forbidden. Furthermore, we envisage this large intermediate scale having its origin in dimension four terms in the superpotential.

b. Model II

This model differs from model I in two respects. First, we add two E_6-singlet fields S_1 and S_2 and impose the discrete symmetry under which J , $\bar{J}$ and S_2 fields are odd and all other fields are even. The superpotential for this model has the form

$$P_{II} = P_0 + P'_{II} \qquad (7)$$

where

$$P'_{II} = \lambda_{14} n_0 \bar{n}_{0,J} S_2 + \lambda_{15} \nu^c \bar{\nu}^c_J S_2 + \lambda_{16} n_0 \bar{n}_{0,H} S_1 + \lambda_{17} \nu^c \bar{\nu}^c_H S_1 \ . \qquad (7a)$$

The second difference from Model I is in the pattern of symmetry breaking which we assume to be as follows:

$$\langle n_{0,H} \rangle = \langle \bar{n}_{0,H} \rangle = v_6$$

and

$$\langle \nu^c_J \rangle = \langle \nu^c_{\bar{J}} \rangle = V_{BL} \; . \tag{8}$$

As a result, N , N^c and n_0 disappear from the low energy spectrum, leaving a 3×3 mass matrix of the following type (assuming $\beta_1 \ll \lambda_{15}$)

$$\begin{array}{c} \\ \nu \\ \nu^c \\ S_2 \end{array} \begin{array}{c} \begin{array}{ccc} \nu & \nu^c & S_2 \end{array} \\ \begin{pmatrix} 0 & m_D & 0 \\ m_D & 0 & V_{BL} \\ 0 & V_{BL} & 0 \end{pmatrix} \end{array} \; . \tag{9}$$

This gives a massless Majorana neutrino:

$$\nu_{phys} \approx \nu - \frac{m_D}{V_{BL}} S_2 \tag{10}$$

and a heavy Dirac neutrino with mass $\approx V_{BL}$. In this approximation, the amplitude for neutrinoless double beta-decay vanishes identically. However, supersymmetry breaking could induce a Majorana mass for S_2 , m_S leading to a small Majorana mass for the neutrino, $m_\nu \approx m_S \cdot (m_D/V_{BL})^2$. For $m_D \approx 1$ MeV for the first generation, $V_{BL} \gtrsim 1$ TeV , a value of $m_S \approx 10^2$–10^3 GeV would lead to $m_{\nu_e} \lesssim .1$ to 1 eV , which is consistent with all observations. Again, this mechanism can be trivially extended to higher generations by adding two E_6-singlets per generation. It may be noted that this mechanism is similar to that discussed in ref. 8.

It is also worth pointing out that one can construct a variant of this model without S_1 , with the same result for neutrino masses. The only new feature is that n_0 remains massless. However, since it does not couple to any light gauge bosons, it remains invisible. This will also not contribute to the expansion of the universe at the nucleosynthesis epoch.

4. In the previous section, we have assumed certain light Higgs multiplets to obtain realistic neutrino masses . We have to show that this can happen without enlarging the gauge group to an unacceptable level. The procedure for deciding this has been outlined in ref. 12. One has to make sure that, under the discrete group $\mathcal{H} \subset E_6$ which is a subgroup of the discrete group in the Calabi-Yau manifold, K , the light fields must remain invariant, i.e.

$$\psi(gx) = U_g\psi(x) \equiv \psi'(x) \qquad (11)$$

where $x \subset K$ and $U_g \subset \mathcal{H}$. It is known that, in order to leave $SU(3)_c \times SU(2)_L$ as an unbroken subgroup following the "flux-loop" breaking mechanism,[13] the U_g is parametrized by six numbers $(x_1,...x_6)$ (i.e. $U_g = e^{ix_iH_i}$ where H_i belong to the Cartan subalgebra of E_6) with $(x_1,...,x_6) = (-c,c,a,b,c\ o)$. It is then straightforward to check that for ν^c and n_o components to remain light, we must have $a = -c$ and $b = 3c$. The unbroken low energy group after flux breaking mechanism is then given by $SU(3)_c \times SU(2)_L \times SU(2) \times U(1) \times U(1)$, if the discrete group $\mathcal{H}$ is Z_n , $n \geq 7$, and $SU(3)_c \times SU(3)_L \times SU(3)_R$, if $\mathcal{H} = Z_3$. In table I, we list the low energy groups for other discrete symmetries. An important point to note is that the low energy electroweak gauge group after the intermediate scale is phenomenologically acceptable in all cases except the ones corresponding to Z_2 , Z_4 and Z_5 symmetry, where SU(5) or SU(6) gauge group survives below the Planck scale. Moreover, the nature of these groups is important in determining the couplings of any extra light Z-boson (as in model II) at low energies.

The next question to ask is where do the light Higgs doublets that break $SU(2)_L \times U(1)_Y$ symmetry and give mass to fermions come from? This depends on the discrete symmetry in question. If the discrete symmetry is Z_3, we find that two $SU(2)_L$ doublets (N,E^-) and (E^+,N^c) from the $\{27\} + \{\overline{27}\}$ Higgs multiplets remain light and can, therefore, serve as light doublets that break $SU(2)_L \times U(1)_Y$ symmetry. These models are similar to the ones discussed by del Aguila et al. (ref. 6). In this case, the constraints of $\sin^2\theta_W$ require that $V_{BL} \approx V_6 \approx 10^{11}$ GeV.

On the other hand, in other cases, where no light $SU(2)_L$-doublet survives from $\{27\} + \{\overline{27}\}$ pair we propose that they come from matter multiplets; more specifically we have in mind the two doublets (per generation) (N,E^-) and (N^c,E^+). One can assign vev's to $\tilde{N}$ and $\tilde{N}^c$ to break $SU(2)_L \times U(1)$. From eqn. (4a), we see that the λ_2, λ_3, λ_9 and λ_{11} terms can then lead to fermion masses. In this class of models, where light Higgs doublets (E,E^c) arise as part of the matter multiplets, we assume as before that both axial B-L and vector B-L symmetry are broken at an intermediate scale. We then have to tune $\lambda_{10} \simeq 0$ to keep the E and E^c light. This is a situation we would like to avoid and our hope is that some day, a discrete symmetry will emerge from an appropriate Calabi-Yau space that will justify the vanishing of λ_{10}.

5. We now discuss baryon number violation in these models. A typical diagram that makes a dominant contribution to proton decay is shown in fig. 1 and we estimate the $\Delta B = 1$ amplitude

$$A_{\Delta B \approx 1} \simeq \frac{\lambda_1 \lambda_4 m_{\tilde{g}}}{M_D M_{Sq}^2} \left(\frac{\alpha_s}{4\pi}\right) . \qquad (12)$$

where $m_{\tilde{g}}$, M_{sq} and M_D represent the gluino, squark and D-quark masses respectively and λ_i are coupling constants in eqn. (4). Choosing $M_D \simeq 10^{12}$ GeV, $m_{\tilde{g}} \approx 10$ GeV, we find, $A_{\Delta B \approx 1} \approx 10^{-19} \lambda_1 \lambda_4$ GeV^{-2}. Since the couplings

λ_i in our superpotential eqn. (4a) are related at the Planck scale, it is reasonable to expect them to be $\approx 10^{-5}$–10^{-6} . Choosing $\lambda_1 \approx \lambda_4 \simeq 10^{-6}$, we find $A_{\Delta B \approx 1} \simeq 10^{-31}$ GeV^{-2} , which is consistent with present experiments. We expect the photino to be heavier than the proton so that proton decay via photino emission is avoided.

Turning now to $\Delta B = 2$ transitions, we first note that it requires $\langle \nu^c_H \rangle$, or $\langle \nu^c_\psi \rangle \neq 0$. Therefore, in the two models for neutrino mass that we have presented here, since $\langle \nu^c_H \rangle = 0 = \langle \nu^c_\phi \rangle$, $\Delta B = 2$ transition is forbidden. On the other hand, in our model II (as well as in other models discussed in literature[9]), one might expect $\langle \nu^c_H \rangle = V^H_{BL} \neq 0$ or $\langle \nu^c_\phi \rangle \neq 0$. In such a case, a non-zero $\Delta B = 2$ amplitude arises from the diagram[14] in Fig. 2. Its magnitude is

$$A_{\Delta B = 2} \simeq \left(\frac{\lambda_5 \lambda^H_7 V^a_{BL}}{M_D}\right)^2 \frac{4\pi\alpha_s}{M^4_{sq} m_{\tilde{g}}} , \quad a = H \text{ or } \phi . \qquad (13)$$

The corresponding n–$\bar{n}$ mixing strength is given by $\delta m_{n-\bar{n}} \simeq A_{\Delta B=2} |\psi(0)|^4$. In our model II, we prefer $V^H_{BL} \ll M_D$ in order not to spoil the neutrino mass results. For instance, if we choose, $V^H_{BL} \lesssim 10^6$ GeV , using $|\psi(0)|^4 \simeq 10^{-3}$ GeV^6 , we find, $\delta m_{n-\bar{n}} \simeq 10^{-25} \lambda_5^2 \lambda_7^2$ GeV , which for $\lambda_5 \approx \lambda_7 \approx 10^{-2}$ can lead to mixing times, $\tau_{n-\bar{n}} \simeq 6 \times 10^9$ sec. This may be barely accessible in future experiments with intense cold neutron beams. In the type II models of neutrino masses,[9] $V^H_{BL} \approx 10^{12}$ GeV . In such models, smaller values of λ ($\sim 10^{-5}$) can also lead to observable n–$\bar{n}$ oscillation. Finally, it is possible that all three fields in the matter field self coupling in eqn. (4) do not belong to the same generation, thereby weakening the constraints on λ_1 and λ_4 . This may improve the situation with respect to n–$\bar{n}$ oscillation.

In conclusion, we have outlined scenarios for understanding small neutrino masses in realistic superstring models. In these models, catastrophic proton decays are avoided. Furthermore, under certain circumstances, they may lead to barely observable neutron-antineutron oscillation.

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Figure Captions

Fig. 1. Box diagram for $\Delta(B-L) = 0$ decay mode of the proton.

Fig. 2. Tree diagram for $\Delta B = 2$ transitions such as N-$\bar{N}$ oscillation.

Discrete Symmetry	Gauge Groups Below Planck Scale (prior to intermediate scale breaking)
Z_2	$SU(6) \times SU(2)_L$
Z_3	$SU(3)_c \times SU(3)_L \times SU(3)_R$
Z_4	$SU(5) \times SU(2)_L \times U(1)$
Z_5	$SU(5) \times SU(2)_N \times U(1)$
Z_6	$SU(3)_c \times U(1) \times SU(2)_L \times SU(3)_R$
Z_n, $n \geq 7$	$SU(3)_c \times SU(2)_L \times SU(2)_N \times U(1) \times U(1)$

Table I: Low energy gauge groups below Planck scale for different choice of discrete symmetries.

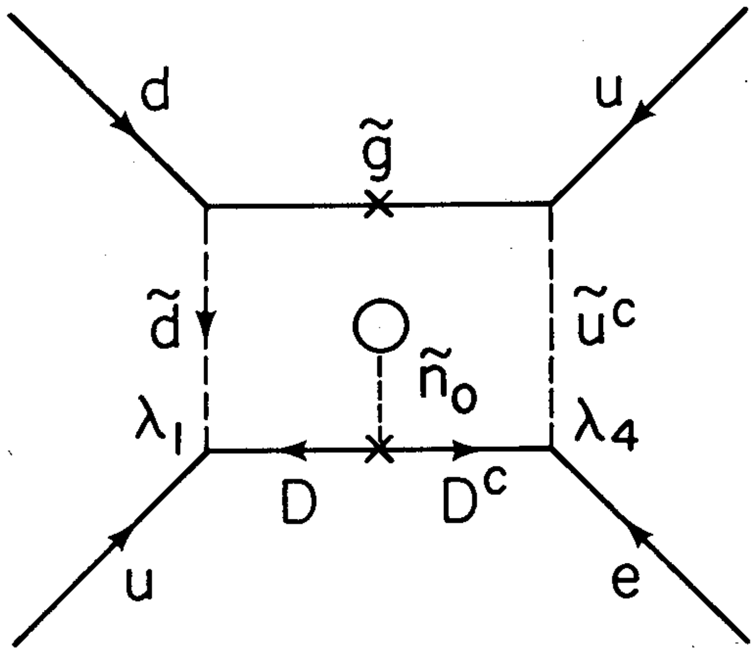

Fig. 1

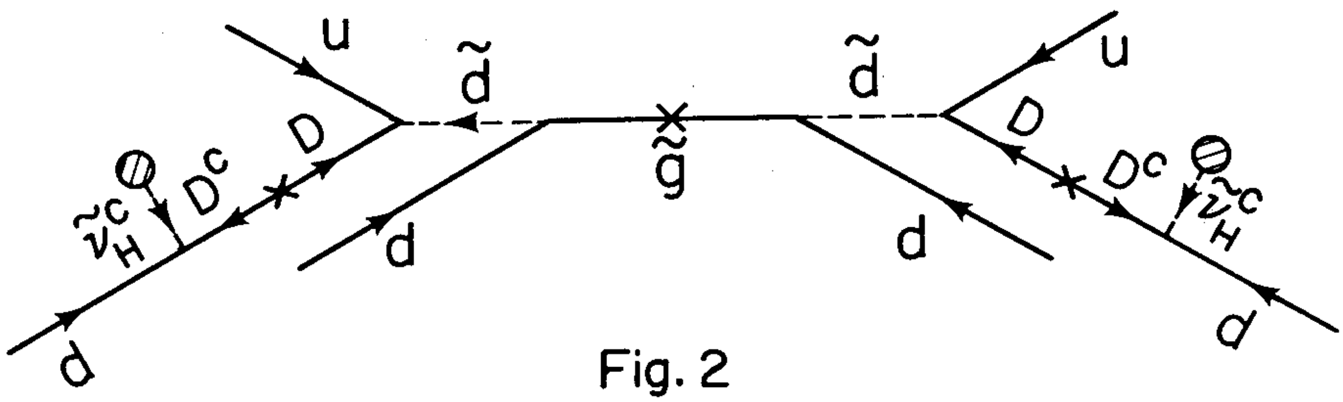

Fig. 2