Ref.TH.3687-CERN

NATURALLY LIGHT DIRAC NEUTRINOS IN GAUGE THEORIES

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ABSTRACT

Motivated by recent experimental results, we discuss the possibility of generating small Dirac neutrino masses in a natural way. We give a mechanism based on symmetry arguments which works in the $SU(2) \times U(1)$ and $SU(5)$ models, but fails in left-right symmetric theories. We conclude that in these models Dirac masses can only be made small by a fine tuning.

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1. - Some three years ago, claims of a non-zero neutrino mass of the order of several eV^{1),2)} led to many theoretical models^{3),4)} for such small masses. The general attitude emerged that only Majorana masses could provide a rationale for their smallness. This point was further corroborated by preliminary evidence for a Majorana electron neutrino mass from neutrinoless double beta decay $(\beta\beta)_0$ ³⁾. Last, but not least, conflicting analyses^{5),6)} of $(\beta\beta)_0$ results were reconciled⁷⁾ within a minimal extension of the standard model^{8),7)}.

Since then, the experimental perspective has changed enough to call these developments into question. While the ITEP result, $14 \text{ eV} < m_{\nu_e} < 46 \text{ eV}$, has remained unchanged in a repeated experiment^{9),*)}, new $(\beta\beta)_0$ experiments^{10),11)} now give $m_{\nu_e} < 15 \text{ eV}$ ^{**)}. While the uncertainties are clearly too big to allow for unambiguous conclusions, the above results, taken at face value, imply a Dirac mass for (at least) the electron neutrino of about 30 eV, and no Majorana mass.

In order to obtain such a small Dirac mass in the standard model, it would be enough to introduce the "right-handed" neutrino and adjust the corresponding Yukawa coupling. This is considered "unnatural" and has been rejected right at the beginning in favour of Majorana masses.

In this paper, we address the question of obtaining naturally a small Dirac mass. We show that - contrary to the old wisdom - there is a positive answer in the $SU(2)_L \otimes U(1)$ and $SU(5)$ models. In order to have naturally small Dirac masses, extra symmetries (to be broken spontaneously) are clearly needed, and along with them, also new fields^{***)}. Nevertheless, the present models are essentially as simple as the original ones. The extra symmetry leads to controllable^{****)} neutrino masses, which in our model are unrelated to the corresponding charged lepton masses. Actually, it will turn out that all neutrinos should weigh almost the same, and, in all cases, below 100 eV.

*) However, under the assumption of a particular model for valine molecular excitations.

**) The actual value of this bound depends on the specific experiment and theoretical interpretation. For example, the geochemical experiment yields $m_{\nu_e} < 6 \text{ eV}$ ¹⁰⁾.

***) This is true for a significantly better understanding of many phenomena; simple book-keeping arguments (fewest fields) should be mistrusted.

****) E.g., no infinite renormalization.

We note that although discrete symmetries are sufficient for our purpose, we prefer continuous (global) $U(1)$ symmetries in order to circumvent the "domain wall problem"¹²⁾. The Goldstone boson we then encounter will turn out to be invisible.

2. - We enlarge the standard $SU(2) \times U(1)$ model by introducing, together with the "right-handed" neutrino, D_R , a pair of Weyl spinors S_L, S_R for each fermion family, which are singlets under the gauge group, but with lepton numbers as the usual neutrinos or charged leptons. Take first a single family only. The most general lepton number-conserving Yukawa Lagrangian is (in the lepton sector):

$$\begin{aligned} \mathcal{L}'_Y = & h_1 \bar{L} \tilde{\varphi} e_R + h'_1 \bar{L} \varphi D_R + h_2 \bar{L} \varphi S_R + \\ & h'_2 \Phi \bar{S}_L S_R + h_3 \Phi \bar{S}_L D_R + M \bar{S}_L S_R + M' \bar{S}_L D_R + h.c. \end{aligned} \quad (1)$$

where $L^T \equiv (D_L, e_L^-)$ and φ are the usual lepton and Higgs doublets, whereas Φ is a new gauge scalar singlet with zero lepton number. Finally, M and M' are arbitrary bare masses. To have naturally light neutrinos, we proceed according to the previously outlined method and impose a symmetry in order to constrain the otherwise largely arbitrary neutrino mass matrix arising from $\mathcal{L}'_Y$. We consider the following symmetry:

$$U(1)_D : D_R \longrightarrow e^{i\theta} D_R, \quad \Phi \longrightarrow e^{i\theta} \Phi \quad (2)$$

with all the other fields invariant. Equation (2) forces $h'_1 = h'_2 = M' = 0$. Observe that the presence of Φ is essential in order to have $h_3 \neq 0$.

The Lagrangian (1) now becomes

$$\mathcal{L} = h_1 \bar{L} \tilde{\varphi} e_R + h_2 \bar{L} \varphi S_R + h_3 \Phi \bar{S}_L S_R + M \bar{S}_L S_R + h.c. \quad (3)$$

When φ and Φ develop vacuum expectation values, $\langle \varphi \rangle$ and $\langle \Phi \rangle$, (3) gives rise to the following tree-level mass matrix for the neutral fermions:

	D_L	S_L	D_R	S_R
$\overline{D}_L$	0	0	0	m_1
$\overline{S}_L$	0	0	m_2	M
$\overline{D}_R$	0	m_2	0	0
$\overline{S}_R$	m_1	M	0	0

(4)

with $m_1 = h_2 \langle \varphi \rangle$ and $m_2 = h_3 \langle \Phi \rangle$. Since $\langle \varphi \rangle$ breaks $SU(2) \times U(1)$, $m_1 < G_F^{-\frac{1}{2}}$. On the other hand, M is a gauge group singlet and the "survival hypothesis"¹³⁾ suggests that it is very large. As we shall see, M is also a singlet under $SU(5)$, so it will generally escape to the cut-off of the theory. At the present state of the art, this is the Planck Mass, $M_p \sim 10^{19}$ GeV. Correspondingly, $M^2 \gg m_1 m_2$; m_2 is also a gauge group singlet, but transforms non-trivially under the $U(1)_D$. This might be used to argue that $\langle \Phi \rangle \ll M_p$ (see later); but we will see [Eq. (11)] that we must assume this relation.

As a consequence, the four Weyl spinors pair up into two Dirac spinors which are eigenstates of (4) as follows (in leading orders in small quantities):

$$\nu \simeq D_L + D_R - \frac{m_1}{M} S_L - \frac{m_2}{M} S_R \quad (5a)$$

$$N \simeq S_L + S_R + \frac{m_1}{M} D_L + \frac{m_2}{M} D_R \quad (5b)$$

with tree-level masses

$$m_\nu \simeq \frac{m_1 m_2}{M} \simeq \frac{h_2 h_3 \langle \varphi \rangle \langle \Phi \rangle}{M} \quad (6a)$$

$$m_N \simeq M \quad (6b)$$

At this point, a comment is very important. For the first family, Eq. (6a) should yield $m_{\nu_e}/m_e \approx 10^{-6} \dots 10^{-5}$, which is such a tiny ratio that it might well be completely upset by radiative corrections. In order to make sure that this is actually not the case, we consider the leading one-loop contribution to Eq. (6a), which comes from the diagram in the Figure. It yields non-vanishing $\overline{D}_L D_R$, $\overline{D}_R D_L$ entries in the neutrino mass matrix of the order:

$$\delta m \approx \frac{\lambda}{4\pi} \cdot \frac{m_1 m_2}{M} \quad (7)$$

Since λ is a quartic Higgs self-coupling, the validity of perturbation theory implies $\delta m < m_1 m_2 / M$, and so Eqs. (6a) and (7) are practically unaffected by this contribution^{*)}.

We now turn our attention to the coupling of ϕ to the fermions. Writing $\phi = 1/\sqrt{2} (\langle \phi \rangle + \rho + i\chi)$, χ is the Goldstone boson of the $U(1)_D$ symmetry, the Diron^{**)}. At tree level, ρ and χ couple only to neutrinos:

$$\begin{aligned} \mathcal{L}_{\rho, \nu N} \approx & \frac{h_3}{2\sqrt{2}} \rho \left[(\bar{N}\nu + \bar{\nu}N) + (\bar{N}\gamma_5\nu - \bar{\nu}\gamma_5 N) \right] + \\ & + \frac{h_3}{\sqrt{2}} \cdot \frac{m_2}{M} \rho \bar{N}N - \frac{h_3}{\sqrt{2}} \cdot \frac{m_1}{M} \rho \bar{\nu}\nu \end{aligned} \quad (8)$$

$$\begin{aligned} \mathcal{L}_{\chi, \nu N} \approx & i \frac{h_3}{2\sqrt{2}} \chi \left[(\bar{N}\nu - \bar{\nu}N) + (\bar{N}\gamma_5\nu + \bar{\nu}\gamma_5 N) \right] + \\ & + i \frac{h_3}{\sqrt{2}} \cdot \frac{m_2}{M} \chi \bar{N}\gamma_5 N - i \frac{h_3}{\sqrt{2}} \cdot \frac{m_1}{M} \chi \bar{\nu}\gamma_5 \nu \end{aligned} \quad (9)$$

in agreement with the general theorem that a Goldstone boson couples only with a γ_5 to diagonal fermions, but also without γ_5 to off-diagonal ones¹⁵⁾. As far as the Diron couplings to charged fermions are concerned, the situation is completely similar to that considered in Ref. 16), and thus the Diron turns out to be invisible.

^{*)} An analogous situation arises in the case of small Majorana masses [see Refs. 3) and 4)], as discussed in Refs. 14).

^{**)} This morpheme comes from a "normalized" combination of "D right" - the only fermion with non-trivial $U(1)_D$ quantum numbers - and "Dirac", since the breakdown of $U(1)_D$ yields Dirac masses.

3. - So far, we have considered only one generation of leptons. For N generations, we have to introduce N pairs S_{Li} , S_{Ri} , along with N D_{Ri} , $i = 1, \dots, N$. In Eq. (3), h_1, h_2, h_3 and M are replaced by $N \times N$ matrices, which we denote by the same letters, and similarly m_1, m_2 and M in (4). The phenomenology of such a situation has been considered in Ref. 17) and we just repeat some important facts. The low-energy phenomenology depends on m_1/M and m_2/M . Since we expect the elements of these matrices to be of the previously discussed orders of magnitude (i.e., $m_1 < G_F^{-1/2}$, $M \gg m_1, m_2$), the qualitative results remain unchanged: there are N naturally light Dirac neutrinos (mainly D_{Li} and D_{Ri}) and N heavy ones (mainly S_{Li} and S_{Ri}). For the small masses, Eq. (6a) generalizes, after diagonalization, to

$$m_{\nu_i} \simeq K_i \frac{\langle \psi \rangle \langle \bar{\Phi} \rangle}{M} \quad (10)$$

where the K_i are linear combinations of products of elements of h_2 and h_3 . It seems natural to take all these elements to be $O(1)$. On the other hand, the Yukawa couplings of charged fermions are all much smaller and show a hierarchical pattern.

By rotating S_{Li} and S_{Ri} ^{*)}, we can choose M to be diagonal. Then, two situations can occur: i) m_1 and m_2 can be diagonalized by only rotating D_{Li} and D_{Ri} ; ii) m_1 and m_2 cannot be diagonalized so simply. Common to both possibilities is the violation of separate lepton numbers - total lepton number is obviously conserved - but in the latter case, the neutrinos produced in β decay are not orthogonal to the ones produced in muon decay; present experimental limits are way above the expected numbers. These considerations also apply to the case of Majorana neutrino masses ^{3), 4)}. Similarly, $\mu \rightarrow e\gamma$ might occur, but from Ref. 17) we again deduce that the present branching ratio is well below the experimental bound. Analogously, small deviations from Cabibbo universality, expected because the light neutrinos are not "fully" in doublets, are negligible. Flavour-changing Higgs-induced processes are also very small, since either the couplings are very tiny, or alternatively the Higgs mass is quite large [see Eqs. (8) and (9)].

These considerations make all neutrino species practically stable in the present model, so the well-known cosmological bound on the masses of long-lived neutrinos ^{**)}:

$$\sum_i m_{\nu_i} \lesssim 100 \text{ eV} \quad (11)$$

^{*)} E.g., by replacing the S_{Li} by a suitable linear combination $S'_{Li} = \sum_j U_{ij} S_{Lj}$, etc.

^{**)} This bound was first derived in Ref. 18) in a more restrictive form. An up-to-date discussion is contained in Ref. 19).

has to be satisfied. We consider this an important restriction: detection of ν_μ or ν_τ with masses larger than ≈ 100 eV would rule out the present model. Now, requiring consistency of (10) and (11), we get, taking $M \sim M_p$, $k_1 \langle \Phi \rangle \lesssim 10^{10}$ GeV. As we said, the $k_i \sim h_2 h_3$ are unknown. If we take them between $\approx 0(1)$ and $\approx 10^{-6}$ [which we consider a reasonable range: $0(1)$ corresponds to a "natural" value and $h_{2,3} \sim 10^{-3}$ to a typical Yukawa coupling], then $\langle \Phi \rangle \ll M$, consistent with our previous assumption^{*)}. Clearly, if $m_{\nu_e} \approx 30$ eV, the above bounds are almost saturated; note that then the neutrino masses cannot follow the same hierarchical pattern as the charged lepton masses.

The reader should appreciate the fact that our result has been obtained without any fine tuning or other exotic assumptions, but rather from symmetry and naturalness arguments along with a cosmological constraint.

4. - The above considerations can be straightforwardly extended to the grand unified model $SU(5)$ ²²⁾, where there is no "right-handed" neutrino. All one needs to do is to add D_{Ri} , S_{Li} , S_{Ri} and the Higgs Φ as $SU(5)$ singlets. Correspondingly, the results are the same as in the $SU(2)_L \otimes U(1)$ model. Obviously, the new particles do not alter the renormalization group analysis²³⁾. On the other hand, in left-right symmetric models²⁴⁾, the neutrinos have necessarily a Dirac mass proportional to $G_F^{-1/2}$. For instance, in $SO(10)$ ²⁵⁾, $m_\nu = m_u$. Under these circumstances, the only way to get a very small eigenvalue in Eq. (4) is by a fine tuning of m_1 , m_2 , M and the $\bar{D}_L D_R$, $\bar{D}_R D_L$ entries, which are typically ≈ 1 MeV for the first generation. If $m_1 \approx 10^2$ GeV, $M \approx 10^{19}$ GeV, then $m_2 \approx 10^{14}$ GeV in order for the fine tuning to be at all possible. As a consequence, parity gets broken at a very high energy. We have not been able to construct an even marginally appealing model that avoids such a fine tuning. However, we notice that in left-right symmetric models, no fine tuning is required if one wishes to have strictly massless neutrinos¹⁷⁾.

ACKNOWLEDGEMENTS

We thank R. Barbieri and S. Pakvasa for having interested us in the question considered in the present work, as well as E. Bellotti, E. Fiorini, B. Pontecorvo and J. Primack for stimulating discussions. We also thank J. Prentki for a critical reading of the manuscript. The work of one of us (M.R.) is supported by the INFN.

^{*)} It could be that in a more complete model, the present $U(1)_D$ can be linked to the Peccei-Quinn $U(1)$ ²⁰⁾. We note with satisfaction that in this case, the above range of $\langle \Phi \rangle$ contains the presently most acceptable breaking scale for the Peccei-Quinn symmetry²¹⁾.

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FIGURE CAPTION

Figure: Leading one-loop contribution to Eq. (6a).