DEUTSCHES ELEKTRONEN-SYNCHROTRON DESY

DESY 86-112
September 1986

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ISSN 0418-9833

NOTKESTRASSE 85 · 2 HAMBURG 52

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Abstract

We consider the new Yukawa couplings involving heavy matter E_6 fields predicted in the framework of Superstring Theories as a source of mass and magnetic moment for the neutrino. Given the experimental bound $m_{\nu_e} < 46$ eV we derive bounds on the neutrino magnetic moment thus generated. Finally, we produce a scenario where the induced magnetic moment has the correct magnitude ($\sim 10^{-11}\ \mu_B$) to explain an alleged depletion of solar neutrino flux during periods of maximum solar activity.

* On leave of absence from Departament de Física de la Universitat Autònoma de Barcelona.

The neutrino is probably the most puzzling and intriguing object in Particle Physics. It is, however, one of the most abundant particles roaming around us. Its mystery stems from the very tiny (if not zero) values of its mass and static parameters. Yet the properties of the neutrino have a profound influence in a number of physical contexts. The neutrino mass, for instance, may notably affect the expansion rate of the Universe and/or explain the large amount of dark matter in the Universe. Thus a major issue in Particle Physics is to determine and understand the values of the neutrino masses. If we endow the standard electron neutrino of the $SU(2)_L \times U(1)$ Weinberg-Salam model with a right handed degree of freedom, then we must simultaneously incorporate in the model a extremely small ($< 10^{-10}$) Yukawa coupling λ_ν related to its mass by $\lambda_\nu \propto G_F^{-1/2} m_\nu$. The unnaturally tiny value of λ_ν makes the picture quite unsatisfactory. In the more ambitious framework of Grand Unified Theories the possibility of fermion number violating Majorana mass terms has been contemplated. In this context a right handed neutrino is easily available so that a bare Majorana mass term as well as a - Higgs induced - Dirac mass term can be suitably combined so as to produce a very light neutrino together with an unobservable very heavy one. This is the well-known "seesaw mechanism" [(1)].

A nonzero neutrino mass makes neutrino oscillations possible and in the astrophysical domain, for example, neutrino oscillations in solar matter [(2-3)] may play a fundamental rôle in explaining the long standing solar neutrino puzzle [(4)]. A potentially important static property of the neutrino is its magnetic moment, μ_ν. For Majorana neutrinos it is obvious that $\mu_\nu = 0$ but for Dirac massive neutrinos a nonvanishing magnetic moment can be of significant relevance in plenty of places in Particle Physics and Astrophysics. In fact, a nonzero value for μ_ν guarantees the existence of magnetic interactions of the neutrinos. The precession of the spin of the neutrino in the stellar magnetic fields has a considerable effect on its effective weak cross-section [5].

The electromagnetic properties of the neutrinos have been repeatedly surveyed in the literature [5-9]. The value of μ_ν is very small in the Standard Model endowed with a right handed neutrino. One finds [6,7]

$$\mu_\nu \simeq 3.2 \times 10^{-19} \left(\frac{m_\nu}{1\ eV}\right) \mu_B \quad (1)$$

where $\mu_B \equiv e\hbar/2m_e c$. For the electron neutrino, e.g., we have $\mu_\nu < 1.5 \times 10^{-17}\ \mu_B$. The contribution from supersymmetrical particles is even smaller [8]

$$\mu_\nu \lesssim 3.5 \times 10^{-20} \left(\frac{m_\nu}{1\ eV}\right) \mu_B \quad (2)$$

Finally, the calculation of μ_ν in the framework of left-right $SU(2)_L x SU(2)_R x U(1)$ electroweak models, gives [9]

$$\mu_\nu < 2.6 \times 10^{-14} \left(\frac{m_\ell}{m_e}\right) \mu_B \quad (3)$$

where m_ℓ is the mass of the charged lepton ℓ associated to ν .

In this paper we evaluate potentially new contributions to the mass and magnetic moment of the neutrino which arise in the context of the recently developed Superstring Theories [10]. In a class of favourite Superstring models [11-13], after dimensional reduction (from 10 to 4 space-time dimensions) via compactification on a Calabi-Yau manifold, one is left with a four-dimensional E_6 gauge group that contains, as a subgroup, the low energy sector of the theory. Each

generation of matter particles is contained in a $\underline{27}$ representation of E_6. Its decomposition into the familiar SO(10) and SU(5) subgroups reads explicitly,

$$(27)_{E_6} = \begin{cases} (\underline{16} + \underline{10} + \underline{1})_{SO(10)} \\ \left[(\underline{10} + \underline{5}^* + \underline{1}) + (\underline{5} + \underline{5}^*) + 1\right]_{SU(5)} \end{cases} \quad (4)$$

We see from this that apart from the well known 15 Standard Model matter particles in each generation, we are led to 12 new extra particles of the fundamental E_6 representation, which are the following

$$\nu^c(1,1),\ H(1,2),\ \bar{H}(1,2),\ D(3,1),\ D^c(3^*,1),\ N(1,1) \quad (5)$$

where in parenthesis we display their $SU(3)_c \times SU(2)_L$ content. So, in particular, we see that there is a right-handed neutrino (which is welcome if we are to have a nonvanishing μ_ν), and an extra color triplet quark.

The most general E_6 invariant superpotential involves the trilinear couplings

$$\alpha_1 Qd^c\bar{H} + \alpha_2 Qu^cH + \alpha_3 Le^cH + \alpha_4 L\nu^cH \quad (6a)$$

$$\beta_1 \bar{H}HN + \beta_2 DD^cN \quad (6b)$$

and

$$\xi_1 DQQ + \xi_2 D^cu^cd^c \quad (6c)$$

or

$$\lambda_1 D^cLQ + \lambda_2 De^cu^c + \lambda_3 D\nu^cd^c \quad (6d)$$

but not both since (6c)+(6d) would lead to a too fast proton decay. [13] In formulae above L and Q are the superfield doublets

$$L \equiv \begin{pmatrix} \nu \\ e \end{pmatrix}_L \quad \text{and} \quad Q \equiv \begin{pmatrix} u \\ d \end{pmatrix}_L$$

Now let us go back to our point. Suppose $\xi_1 = \xi_2 = 0$ and therefore the terms (6d) are allowed. We shall furthermore assume that $\lambda_3 \leq O(10^{-2})$ in order to prevent a radiatively triggered nonzero value for $\langle 0 | \tilde{\nu}^c | 0 \rangle$. In fact, for $\langle 0 | \tilde{\nu}^c | 0 \rangle \neq 0$ the third term in eq. (6d) would generate D-d mixing and lead to conflict with FCNC constraints [14]. The terms in eq. (6a) give masses to quarks and leptons. It is immediately apparent that, unless $\alpha_4 \ll \alpha_1, \alpha_2, \alpha_3$, an unacceptably large neutrino mass will ensue. It would be, therefore, desirable to ban the fourth Yukawa coupling in eq. (6a). The origin of it (as well as of the previous assumption that $\xi_1 = \xi_2 = 0$) could be attributed either to topological reasons [13] - maybe the overlap integrals defining these Yukawa couplings vanish because some of the complex hypersurfaces do not share a common intersection - or to some conventional discrete global symmetries whose fundamental origin remains unclear. Anyway, if we stick to our assumption ($\alpha_4 = \xi_1 = \xi_2 = 0$) a radiatively induced Dirac neutrino mass is still possible. The relevant pieces are the first and third terms in eq. (6d). The one-loop diagrams depicted in Fig. 1 render the following finite amount

$$m_\nu = \frac{3\,\lambda_1 \lambda_3}{32\,\pi^2} \sin 2\chi_d \, m_D \left[\log \frac{m_1^2}{m_2^2} + \frac{m_D^2}{m_D^2 - m_1^2} \log \frac{m_D^2}{m_1^2} - \frac{m_D^2}{m_D^2 - m_2^2} \log \frac{m_D^2}{m_2^2} \right] \qquad (7)$$

where $m_{1,2}$ stand for the d-squark masses * (recall that squarks come in two

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* For the purpose of illustration we particularize our analysis to the first (fermion-sfermion) generation, i.e., the mass in eq. (7) is the electron neutrino mass.

chiralities) and where χ_d is the angle that diagonalizes the d-squark mass matrix in left-right space,

$$\begin{pmatrix} m_{LL}^2 , & m_{LR}^2 \\ m_{RL}^2 , & m_{RR}^2 \end{pmatrix} \qquad (\text{with } m_{LR}^2 = m_{RL}^2) \qquad (8)$$

The above mass matrix is nondiagonal because supersymmetry breaking introduces off-diagonal entries which in the context of supergravity models have the structure [12,13,15]

$$m_{LR}^2 = A\, m_d\, m_o \qquad (9)$$

where A is a model dependent unknown constant, m_d is the d quark mass and m_o is the supersymmetry breaking scale. It is patent in the result, eq. (7), and obvious from the structure of the diagrams in Fig. 1, that the resulting contribution to m_ν is proportional to the exotic D-quark mass. Also, it is worth mentioning that one needs left-right mixing in the d-squark sector in order to get $m_\nu \neq 0$. Now the D-quark mass, as well as the extra Z' boson mass, is generated from the vacuum expectation value of the E_6 singlet field N in eq. (6b). Small mixing in the mass matrix for the Z and Z' neutral vector bosons requires that $\langle 0|N|0 \rangle \gtrsim O(100)$ GeV, so that this is also the natural scale expected for the mass of the D-quark.

Let us evaluate formula (7) in a few presently popular Superstring inspired models. Two of our choices correspond to the so-called hybrid dimensional-transmutation models [16,17]. The other three are taken from Ref. (18). In Table 1 we display the upper bounds on the product $\lambda_1 \lambda_3$ following from the experimental requirement $m_{\nu_e} < 46$ eV for the various models selected. All generated masses are

$$m_\nu \sim O(10^{-1})\,\lambda_1\lambda_3 \quad MeV \tag{10}$$

and, in fact, a typical choice $\lambda_1 \sim \lambda_3 \sim 10^{-2}$ (quite natural if we compare with the $< O(10^{-10})$ Yukawa couplings needed for tree level m_ν generation) suffices to rightly account for the neutrino mass without any concern about FCNC induced by D-d mixing.

Our mechanism for m_ν generation is quite simple and relies on some unknown topological property or discrete symmetry forbidding the α_4 Yukawa term in eq. (6a), something which has become a common article of faith in many models [13,16-18].

Within the same framework we consider the contribution to the magnetic moment μ_ν of the neutrino induced by the first and third terms in the superpotential (6d). The corresponding Feynman diagrams are shown in Fig. 2. A calculation of these diagrams leads to the following formula,

$$\mu_\nu = \frac{\lambda_1\lambda_3}{8\pi^2}\sin 2\chi_d \sum_{i=1,2}(-1)^i \frac{m_e m_D}{m_i^2 - m_D^2}\left[1 - \frac{m_i^2}{m_i^2 - m_D^2}\log\frac{m_i^2}{m_D^2}\right]\mu_B \tag{11}$$

where the notation is that of eq. (7). Again to provide a scenario we shall evaluate formula (11) in the same Superstring models used to evaluate formula (7). Of course, the bounds displayed in Table 1 for the product $\lambda_1\lambda_3$ are to be used here to obtain the corresponding ones for the neutrino magnetic moment. These upper bounds are exhibited in the last column of Table 1. We see, after comparing them with eqs. (1)-(3), that in general they are bigger than the

standard model and supersymmetrical standard model values, while they are smaller than the value obtained from left-right electroweak models.

The bounds on $\lambda_1\lambda_3$ and, consequently, on μ_ν displayed in Table 1 are absolutely unavoidable if one has $\alpha_4 = 0$ in eq. (6a) to start with. However, it is possible to device a situation where the aforementioned bounds are relaxed. Indeed, suppose that the Yukawa coupling α_4 is not forbidden. This is something that might be natural to think of, since after all the topological origin of these Yukawa couplings is the same. The tree level mass of the neutrino, however, is obviously unacceptable (unless $\alpha_4 \lesssim O(10^{-10})$, which is contrived). An enhanced value for the product $\lambda_1\lambda_3$ could then provide for a radiatively induced mass that exactly conspires with the tree level term just to give the neutrino its tiny physical mass. Although this is technically legal it demands a fine tuning of some two to three decimal places for a neutrino mass ~ 10 eV, which is very mild if compared to GUT's fine tunings. It is, however, the price one has to pay for not introducing from the start, and by hand, an unnaturally tiny α_4 Yukawa coupling. We consider this possibility worth exploring, specially having in mind that it could lead, as we shall see, to experimentally testable consequences in a field otherwise rather meagre in phenomenological implications.

Of course, we must keep $\lambda_3 \sim O(10^{-2})$ in order not to trigger D-d mixing by a radiatively driven $\langle \tilde{\nu}^c \rangle \neq 0$ and therefore $|\lambda_1\lambda_3|$ cannot be too large. Let us choose for the sake of definiteness $\lambda_1 = g_W$ and $\lambda_3 = 3 \times 10^{-2}$ so that $\lambda_1\lambda_3 \cong 2 \times 10^{-2}$.

We have been playing before with some models in the market, but at the present state of the art, one should not stick too literally to the existing models.

One can then get larger values for the magnetic moment of the neutrino and still not contradict any phenomenological constraint. For example, take $m_D = 100$ Gev and the following values

$$m^2_{LL} = (120 \text{ Gev})^2,\ m^2_{RR} = (110 \text{ GeV})^2,\ m^2_{LR} = (3 \text{ GeV})^2 \qquad (12)$$

for the entries in the mass matrix (8). We have chosen m^2_{LR} small so as to conform with the supergravity [15] derived prejudice: $|m^2_{LL} - m^2_{RR}| \gg m^2_{LR}$. Then using the optimized value previously stated for $\lambda_1 \lambda_3$ we get from eq. (11) for the magnetic moment of the neutrino

$$\mu_\nu \simeq 3 \times 10^{-13}\ \mu_B \qquad (13)$$

which is far larger than the results previously obtained in Table 1 when assuming that the α_4 term in the superpotential was absent. It is even an order of magnitude larger than the maximum contribution, eq. (3), expected from left-right electroweak models.

Obviously, the set of parameters used to obtain eq. (13) have no special significance and serve only to illustrate the plausibility of generating a relatively large value of the neutrino magnetic moment from a superstring induced E_6 superpotential. The phenomenological relevance of having a nonzero value for μ_ν was stated already in the beginning of the paper. Here we would like to turn our attention to a specific astrophysical problem for which a substantial magnetic moment of the neutrino is of maximum interest.

Allegedly, the flux of neutrinos from the sun [4] attains its maximum value when the solar activity is less and viceversa. This flux variation proceeds over a

period of eleven years and during this time the neutrino flux changes roughly by a factor of two [19]. A magnetic moment of the electron neutrino of the order of $(10^{-11}\text{-}10^{-10})\mu_B$ would suffice for the magnetic field of the solar convective zone to flip the helicity of the neutrinos by the amount necessary to explain the decrease in the detected flux [20]. This explanation would generate a periodic modulation of the detected flux linked to the solar activity cycle.

As we have seen before (eqs. (1) and (2)) neither the Standard Model nor its supersymmetrical version can account for such a "large" μ_ν. The authors of Ref. (20) claim that if we allow for substantial electron neutrino mixing with a neutrino of a heavier generation (the τ-family or a postulated fourth generation) then the left-right electroweak model can provide for the necessary value for μ_ν. In fact, taking the upper bound (3) and multiplying it by a Cabibbo mixing factor $\sin^2\theta = 1/2$ for $\nu_e - \nu_\tau$ oscillations one finds $\mu_\nu = O(10^{-11})\ \mu_B$ which is of the order required to explain the effect. Notice, however, that in taking the upper bound (3) they assume a mixing factor $\sin 2\phi$ for the left-right charged weak boson states of 0.1 which is almost ruled out by experiment [21].

We would like to point out that a similar result can be accommodated in the superstring framework. Suppose that we also allow for $\nu_e - \nu_\tau$ mixing with $\theta = \frac{\pi}{4}$. Then the scalar particle looping around the diagrams in Fig. 2 will be the b-squark. Consequently, the off diagonal terms in the mass matrix (8) will be m_b/m_d times larger (see eq. (9)). Repeating the calculations which led us to eq. (13) we now obtain

$$\mu_\nu \simeq 2 \times 10^{-11}\ \mu_B \qquad (14)$$

which is about the right size to explain the anticorrelation of the neutrino flux data with the solar activity. Notice that the value in eq. (14) is below the upper limit derived from the neutrino luminosity estimates of white dwarfs [22]. Of course, still larger values could be attained if we tolerate larger mixing angles and/or left-right mass splittings in the matrix (8) or even if we allow for smaller D and/or b-squark masses.

To conclude let us recapitulate our findings.

Starting with a Lagrangian with no mass term for the neutrino we have computed a radiatively induced neutrino mass generated by Yukawa terms in the E_6 invariant superpotential suggested by Superstring Theories. We have bound the two relevant Yukawa couplings by requiring that the radiative corrections do respect the experimental bound on m_ν. The bounds obtained for the product of the two Yukawa couplings are not unnaturally tiny. Individual values of $O(10^{-2})$, e.g., are allowed, in contradistinction to those that would be needed at tree level in order to give the neutrino a mass. These results have been obtained by scanning on a handful of models in the market. We then used the bounds obtained from the ν mass to compute potentially new contributions to the neutrino magnetic moment μ_ν. Although the parameters in the models presently available tend to give low values for μ_ν they are nevertheless larger than the Standard Model one. Moreover, we have shown that still larger values for μ_ν are possible if we allow for a Yukawa term $L\nu^c H$ with a coupling of considerable strength (not much weaker than that of the $Le^c H$ term). To elude a phenomenologically unacceptable neutrino mass we required that radiative corrections restore a sensible value for m_ν. The fine tunning needed is rather soft (at the 10^{-3}-10^{-2} level) if one compares it to other more dramatic ones occurring in

GUT's. On the other hand this is a mere reflection of a well-known difficulty in Superstring modelling. It is not known how to implement in a natural way a small neutrino mass [23]. We have envisaged this possibility mainly because of its phenomenological consequences. It leads to far larger μ_ν values, which can be derived with a phenomenologically sound choice of masses, mixings and couplings. Indeed we have shown that in this framework we can accommodate a μ_ν value that may help explain an alleged depletion of the neutrino flux linked to the solar activity cycle. Probably it is too premature to draw definite conclusions from present solar data but in any event the experimental study of the time variation of the solar neutrino flux and its correlation to the solar activity should be actively pursued on an improved statistical basis in the forthcoming years. This is an important issue that can shed light both on the models of solar structure and on the properties of neutrinos. If confirmed, the need for a substantial μ_ν value becomes more pressing and its natural source might well be found in the Superstring context provided, of course, that the final and natural solution (if existing) to the neutrino mass problem does not radically constrain the μ_ν-inducing terms.

Acknowledgements

We are grateful to R.D. Peccei for bringing the mass problem to our attention, for a critical reading of the manuscript and for fruitful discussions, and to C. Wetterich for useful comments. J.A.G. would like to thank Prof. R.D. Peccei for his kind invitation to spend a month at DESY and we both thank the DESY theory group for its hospitality.

Figure Captions

Fig. 1 Feynman diagrams contributing to the mass of the neutrino. The cross on a fermionic line stands for a mass insertion and a cross on a scalar line stands for left-right mixing.

Fig. 2 Feynman diagrams contributing to the magnetic moment of the neutrino. The notation is like in Fig. 1.

Table Caption

Table 1

Upper bound values for $|\lambda_1 \lambda_3|$ and μ_ν for the various models quoted. Since Table 2 of Ref. (18) does not separately specify left-right scalar quark masses, we have assumed a ~5 % mass splitting (consistent with the splittings given in Table 4 of Ref. (16)). All masses are given in GeV.

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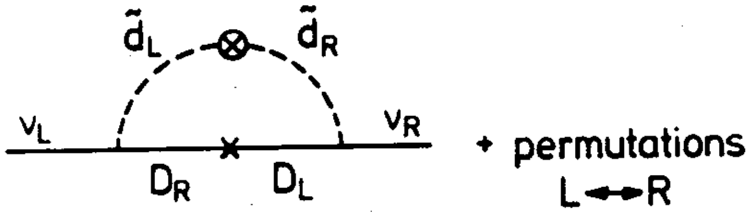
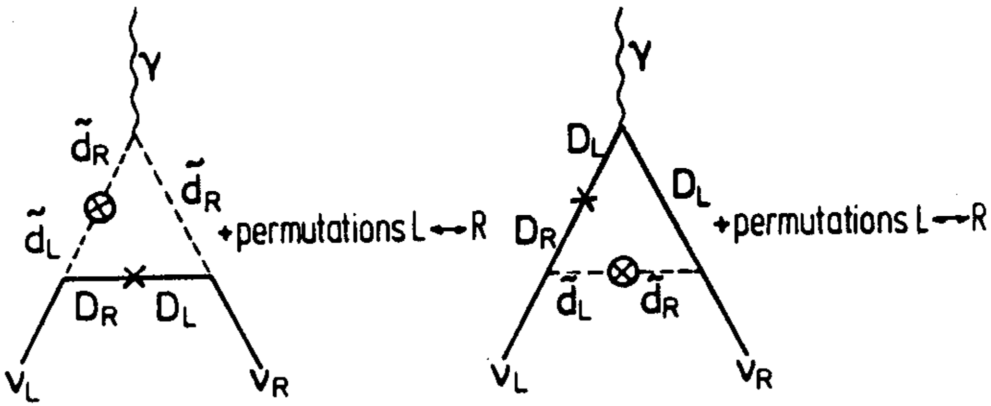
Table I

	m_1	m_2	m_D	A	$\lvert\lambda_1\lambda_3\rvert$	μ_ν (in μ_B units)
Model (a) Ref. 16	1600	1500	530	3	$3x10^{-4}$	$1.5x10^{-17}$
Model (b) Ref. 16	610	560	180	3	$3x10^{-4}$	$1x10^{-16}$
Model (a) Ref. 18	263	250	222	2.6	$2x10^{-4}$	$2x10^{-16}$
Model (e) Ref. 18	315	300	438	-0.86	$5x10^{-4}$	$6.5x10^{-17}$
Model (h) Ref. 18	131	125	206	-2.1	$2x10^{-4}$	$3x10^{-16}$



(Fig.1)



(Fig.2)

ADDED NOTE

After our manuscript was sent to the printer, we became aware of a paper by A. Masiero, D.V. Nanopoulos and A.I. Sanda (Phys. Rev. Lett. 57 (1986) 663) where part of our discussion (the one relative to the neutrino mass problem) is also considered, though from another point of view. Their mechanism for neutrino mass generation uses D-squark scalars instead of D-quark fermions. The numerical contribution, however, is of the same order of magnitude and can be added to ours without modifying our conclusions. On the other hand, their bounds on the Yukawa couplings depend on assumptions on the D-squark mass and on intergenerational mixings, while our Yukawas refer to the same generation and no use is made of the D-squark mass. Both pictures are, therefore, independent and complementary.