A GUT SOLUTION TO THE SOLAR NEUTRINO PROBLEM

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Within the supersymmetric flipped SU(5)×U(1) model, we propose a mechanism for realization of the Voloshin–Vysotsky–Okun solution to the solar neutrino problem by attributing a large magnetic moment to the electron neutrino, as required to explain the solar neutrino data.

One of the long-standing problems in astrophysics is the solar neutrino one [1], namely that the number of energetic neutrinos captured on earth [2] is almost three times smaller than their predicted number emitted by the sun, according to the standard solar model [3]. During the past years, many solutions have been proposed to solve this problem. These include neutrino oscillations in the vacuum [4], resonant neutrino oscillations in solar matter as well as in the earth [5], and the possible existence of suitable weakly-interacting massive particles (called "WIMPs" or "cosmions") [6]. More recently, another interesting suggestion was put forward by Voloshin, Vysotsky and Okun [7]. They showed that, if the electron neutrino had a high magnetic moment of the order of

$$\mu_{\nu_e} \sim (0.3\text{–}1)\times 10^{-10}\mu_B\ , \qquad (1)$$

where $\mu_B = e/2m_e$ is the Bohr magneton, a significant fraction of left-handed neutrinos would precess into sterile right-handed neutrinos under the strong magnetic field in the convective zone of the sun, and thus escape detection. Interestingly enough, this proposal provides a natural explanation of the anticorrelation between the measured neutrino flux and the sun-spot number, as suggested by the neutrino data [8]. This happens because the precessing rate will be high during the period of maximum solar activity, when the magnetic field in the convective zone of the sun will be strong. On the other hand, precession will be very low during the quiet period of solar activity, when the magnetic field is at least one order of magnitude weaker. As a result, a higher capture rate in the earth is expected in this period [7]. Another feature of this mechanism is the predicted biennial variation of the solar neutrino flux [7], for which there also seems to be some evidence. Forthcoming solar neutrino detectors are expected to look for such phenomena.

Note that the value (1) is just below the experimental upper bound set by laboratory experiments from ν–e scattering ($\mu_{\nu_e} \lesssim 1.5\times 10^{-10}\mu_B$) [9] and by the astrophysical consideration of stellar cooling ($t_{\nu_e} \lesssim 1\times 10^{-10}\mu_B$) [10].

Many authors [11] have tried to find models to realize such a large magnetic moment for the neutrino, which is generally negligible in the standard model and most of its extensions [12]. In almost all cases, a considerable deviation from the standard SU(2)×U(1) theory is required. One also has to be careful to satisfy an additional constraint put by cos-

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mology. The constraint here [13] comes from the fact that an upper limit

$$\mu_\nu < 1.5\times10^{-11}\,\mu_B \qquad (2)$$

should not be violated by more than two neutrino species. Otherwise, $\nu_L e\to\nu_R e$ scattering (due to a large magnetic moment), before neutrino decoupling, would double the effective number of neutrino species in the early universe, thus causing an excess of ^{4}He. This cosmological constraint is satisfied if the bound (2) is violated only for μ_{ν_e}, as in (1), but not for μ_{ν_μ}, μ_{ν_τ}, which must be $\ll\mu_{\nu_e}$.

In the present letter, we shall show how to realize the Voloshin–Vysotsky–Okun (VVO) poposal for the solar neutrino problem in a grand unified theory.

Let us first discuss the situation within the conventional SU(5) model. In order to get a large magnetic moment for the ν_e and not for the ν_μ and ν_τ, we need to have antisymmetric (in the generation space) Yukawa couplings. Then, besides the usual matter (super) fields F(**10**), F($\bar{\mathbf{5}}$) and Higgses H(**24**), H(**5**) [H($\bar{\mathbf{5}}$)], we have to introduce an antisymmetric representation of Higgs H(**10**) [H($\overline{\mathbf{10}}$)]. If we also introduce the right-handed neutrino singlet F(**1**), the invariant couplings F($\bar{\mathbf{5}}$)F($\bar{\mathbf{5}}$)H(**10**), F(**1**)F(**10**)H($\overline{\mathbf{10}}$) will provide us with the required Yukawa terms:

$$F(\bar{5})\,F(\bar{5})\,H(10)\to c_{ij}\nu_{iL}e^c_{jR}S\,, \qquad (3)$$

$$F(1)\,F(10)\,H(\overline{10})\to d_{ij}\nu_{iR}e_{jL}\bar{S}\,. \qquad (4)$$

Then one could repeat the analysis made in the already quoted work [11]. However, here we encounter a big problem: the above terms (3) and (4) introduce extra couplings, which violate baryon and/or lepton number. In order not to have any conflict with observation (e.g., rapid proton decay), we have to achieve a natural mass splitting for the H(**10**) fields, by keeping the mass of the singly-charged S particles relatively low, $\lesssim O(10^2\text{–}10^3)$ GeV, while giving at the same time a superheavy mass $\sim O(M_{GUT})$ to the remaining H(**10**) particles. One cannot achieve such a mass splitting in the conventional SU(5) model, since there are no natural mechanisms at hand to do that. This is our main motivation to go to the flipped SU(5)×U(1) model, where we can easily solve this problem, as we will argue next.

The flipped SU(5)×U(1) model [14] has some nice features: it does not need any adjoint (or large self-conjugate) Higgs representation to break the GUT symmetry. Its Higgs sector contains the lowest representations **10** and **5**. For our purposes this is welcome, because it is exactly a **10** representation that we want to exploit. So there is no need to introduce Higgs representations other than those already present in the model. It also solves naturally the doublet–triplet mass splitting in the **5**-plet Higgs. Moreover, one can naturally have a seesaw mechanism for neutrino masses, which will also fit into our programme. The model can, in principle, be derived [15] from superstrings. In particular, it can be derived [16] from the four-dimensional (fermionic) formulation [17] of superstrings [#1]. In the last case, one generally expects no adjoint chiral superfields [19].

The basic SU(5)×U(1) model contains three generations of matter fields, $F_i=(\mathbf{10},1)$, $\bar{f}_i=(\mathbf{5},-3)$ and $\ell^c_i=(\mathbf{1},5)$ with the following particle assignments for the first generation:

$$F=\begin{pmatrix} 0 & d^c & -d^c & d & u \\ -d^c & 0 & d^c & d & u \\ d^c & -d^c & 0 & d & u \\ -d & -d & -d & 0 & \nu^c \\ -u & -u & -u & -\nu^c & 0 \end{pmatrix},$$

$$\bar{f}=\begin{pmatrix} u^c \\ u^c \\ u^c \\ \nu \\ e \end{pmatrix} \qquad (5)$$

and similarly for the other generations. The Higgs representations are the conjugate pairs $H=(\mathbf{10},1)$, $\bar{H}=(\overline{\mathbf{10}},-1)$, $h=(\mathbf{5},-2)$, $\bar{h}=(\bar{\mathbf{5}},2)$. The model contains also four SU(5)×U(1) singlets $\phi_m=(\mathbf{1},0)$, three of which play a role in the neutrino seesaw mechanism, and the other one provides the necessary $h\bar{h}$ mixing. The superpotential of the model is

$$W=\lambda_1^{ij}\,F_iF_jh+\lambda_2^{ij}\,F_i\bar{f}_j\bar{h}+\lambda_3^{ij}\bar{f}_i\ell^c_jh+\lambda_4HHh+\lambda_5\bar{H}\bar{H}\bar{h}$$
$$+\lambda_6^{im}F_i\bar{H}\phi_m+\lambda_7^m h\bar{h}\phi_m+\lambda_8^{mnp}\phi_m\phi_n\phi_p\,. \qquad (6)$$

A linear combination of H and $\bar{H}$ acquires a vacuum

[#1] For a mechanism of supersymmetry breaking in these theories, see ref. [18].

expectation value (VEV), which breaks $SU(5)\times U(1)\rightarrow SU(3)_c\times SU(2)_L\times U(1)$, and a VEV for h and $\bar{h}$ breaks $SU(2)_L\times U(1)_Y\rightarrow U(1)_{em}$. All the above λ-couplings have an important phenomenological role to play [14].

As discussed above, in order to implement our programme and give a magnetic moment to the neutrinos, we need a H(**10**) Higgs field. So we introduce the conjugate pair of Higgses $H'=(\mathbf{10},6)$, $\bar{H}'=(\overline{\mathbf{10}},-6)$, and extend the superpotential (6) to include the terms

$$\lambda_9^{ij}\bar{f}_i\bar{f}_j H' + \lambda_{10}^{ij} F_i \ell_j^c \bar{H}' \ . \tag{7}$$

However, these terms also introduce baryon-number-violating interactions. So, as remarked earlier, in order to avoid disastrous consequences, we have to give large masses to all the baryon-number-violating particles in H', $\bar{H}'$, while keeping the singlets S ($\bar{S}$) in them at a low mass $\lesssim O(10^2\text{–}10^3)$ GeV.

As already said, the flipped $SU(5)\times U(1)$ model can be obtained from the four-dimensional fermionic formulation of strings. A first attempt towards this direction is discussed in ref. [16]. Then, in addition to the massless sector, one also obtains massive states, whose mass scale is naturally $O(M_P)$. In particular, our H', $\bar{H}'$ fields will naturally have masses $O(M_P)$. Now, in order to achieve the mass splitting to which we referred before, we will have to make use of a seesaw mechanism. For that, it is sufficient to have an additional $S'=(\mathbf{1},5)$ [note its quantum numbers similar to those of $\ell^c=(\mathbf{1},5)$] and write two more terms in the superpotential

$$\lambda_{11}\bar{H}\bar{S}'H' + \lambda_{12}HS\bar{H}' \ . \tag{8}$$

These two terms, through the VEV $\langle H\rangle = \langle \bar{H}\rangle \sim O(M_{GUT})$, will provide us with the required see saw mechanism. For $\lambda_{11,12}\sim O(10^{-3}\text{–}10^{-4})$, we will have a mass $\sim O((10^{11})^2/M_P)\sim O(10^3)$ GeV for one combination of S, S', whereas all the other components will naturally have masses $\sim O(M_P)$. So we have achieved the required mass splitting within the 10-plets of Higgses H', $\bar{H}'$.

We are now ready to go on to our point, i.e., to calculate the magnetic moment of the neutrinos. The Yukawa terms (7) give rise to couplings as in (3) and (4) generating four one-loop diagrams shown in fig. 1. The first two, (a) and (b), are the same as in the non-supersymmetric cases considered previously

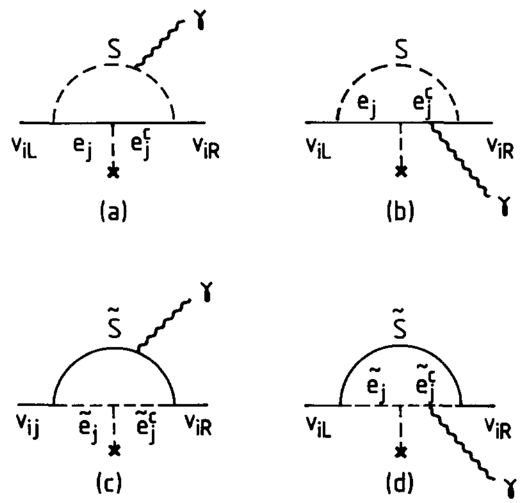

Fig. 1.

[11], whereas the last two, (c) and (d), are their supersymmetric analogues. The contribution to the magnetic moment from (a) and (b) is, as previously [11],

$$\mu_{\nu_i} = \frac{e}{2m_e}\sum_{j\neq i}\frac{d^\dagger_{ij}c_{ji}+c^\dagger_{ij}d_{ji}}{16\pi^2}\,\frac{m_j m_e}{m_S^2}\left(\ln\frac{m_S^2}{m_j^2}-1\right), \tag{9}$$

whereas for the corresponding contribution from the supersymmetric graphs we have

$$\mu_{\nu_i} = \frac{e}{2m_e}\sum_{j\neq i}\frac{c^\dagger_{ij}d_{ji}+d^\dagger_{ji}c_{ji}}{32\pi^2}\,\frac{m_e\tilde{m}_j^2}{\tilde{m}_S^3}\left(\ln\frac{\tilde{m}_S^2}{\tilde{m}_j^2}-3\right). \tag{10}$$

In the above formulae, m_j stands for the charged lepton mass of the jth generation, $\tilde{m}_j$ for its supersymmetric partner mass, m_S is the mass of the scalar S particles and $\tilde{m}_S$ that of their superpartners. In both expressions (9) and (10), the magnetic moment of the ν_i picks up the mass of the jth ($j=i$) charged particles and their superpartners. Therefore ν_e and ν_μ are expected to have a larger magnetic moment. The numerical estimate of the magnetic moments (9) and (10) depends on the value of the parameters involved. By choosing reasonable values for them, e.g., $\tilde{m}_j\sim 10^2$ GeV, $\tilde{m}_S\sim 500$ GeV, $m_S\lesssim 10^3$ GeV, we find that the dominant contribution comes from the supersymmetric expression (10), which will be of the right order (1) for c_{13}, $d_{31}\sim O(10^{-1})$. Furthermore, to satisfy the cosmological constant (2), one has to take the Yukawa couplings c_{23}, d_{32} entering the expression for the ν_μ to be one order of magnitude smaller. We have checked that this range of parameters predict branching ratios for the various flavour-changing processes [20], which lie safely within the present experimental bounds [21].

Finally, a comment on the neutrino masses in the present model. The same graphs as in fig. 1, but without the external photon line, give radiative contributions to the neutrino mass terms. The non-supersymmetric graph is logarithmically divergent, whereas the supersymmetric one gives a contribution

$$m_{\nu_{ij}} = \sum_k \frac{(d^\dagger c + c^\dagger d)_{ij}}{16\pi^2} \frac{\tilde{m}_k^2}{\tilde{m}_S} \left(\ln \frac{\tilde{m}_S^2}{\tilde{m}_k^2} - 1 \right) . \quad (11)$$

This contribution will be added to the off-diagonal terms of the neutrino mass matrix in the seesaw mechanism operative in the present model [1 9]. So the neutrino mass remains naturally small.

In summary, we have made an attempt to find, in a grand unified theory, a way to realize a large magnetic moment for the electron neutrino, necessary to explain the solar neutrino problem. We have shown that this can be naturally achieved in a flipped $SU(5)\times U(1)$ model, for a range of parameters consistent with the experimental observations.

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