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Neutrino Masses and the Gluino Axion Model

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Abstract

We extend the recently proposed gluino axion model to include neutrino masses. We discuss how the canonical seesaw model and the Higgs triplet model may be realized in this framework. In the former case, the heavy singlet neutrinos are contained in superfields which do not have any vacuum expectation value, whereas the gluino axion is contained in one which does. We also construct a specific renormalizable model which realizes the mass scale relationship $M_{SUSY} \sim f_a^2/M_U$, where f_a is the axion decay constant and M_U is a large effective mass parameter.

A new axionic solution[1] to the strong CP problem was recently proposed[2]. Instead of coupling to ordinary matter as in the DFSZ model[3] or to unknown matter as in the KSVZ model[4], this new axion couples to the gluino as well as all other supersymmetric particles. The instanton-induced CP violating phase[5] of quantum chromodynamics is then canceled by the dynamical phase of the gluino mass, as opposed to that of the quarks in the DFSZ model and that of the unknown colored fermions in the KSVZ model. This means that CP violation is absent in the strong-interaction sector and experimental observables, such as the neutron electric dipole moment[6], are subject only to weak-interaction contributions.

What sets the gluino axion model[2] apart from all other previous models is its identification of the Peccei-Quinn global symmetry $U(1)_{PQ}$ with the $U(1)_R$ symmetry of superfield transformations. Under $U(1)_R$, the scalar components of a chiral superfield transform as $\phi \to e^{i\theta R}\phi$, whereas the fermionic components transform as $\psi \to e^{i\theta(R-1)}\psi$. In the Minimal Supersymmetric Standard Model (MSSM), the quark and lepton superfields $\hat{Q}$, $\hat{u}^c$, $\hat{d}^c$, $\hat{L}$, $\hat{e}^c$ have $R = +1$ whereas the Higgs superfields $\hat{H}_u$, $\hat{H}_d$ have $R = 0$. The superpotential

$$\hat{W} = \mu \hat{H}_u \hat{H}_d + h_u \hat{H}_u \hat{Q} \hat{u}^c + h_d \hat{H}_d \hat{Q} \hat{d}^c + h_e \hat{H}_d \hat{L} \hat{e}^c \tag{1}$$

has $R = +2$ except for the μ term (which has $R = 0$). Hence the resulting Lagrangian breaks $U(1)_R$ explicitly, leaving only a discrete remnant, i.e. the usual R parity: $R = (-1)^{3B+L+2J}$. The gluino axion model replaces μ with a singlet composite superfield of $R = +2$ so that the resulting supersymmetric Lagrangian is invariant under $U(1)_R$. It also requires all supersymmetry breaking terms to be invariant under $U(1)_R$, the spontaneous breaking of which then produces the axion and solves the strong CP problem.

In the MSSM, neutrinos are massless. However, in view of the recent experimental evidence for neutrino oscillations, it is desirable to incorporate into any realistic model naturally small Majorana neutrino masses[7, 8]. In the following we will discuss how the canonical seesaw model[9] and the Higgs triplet model[10] may be realized in the framework of the gluino axion model. In the case of the seesaw model, there are in fact proposals[11] that the axion scale is the same as that of the singlet neutrino masses.

Consider first the Higgs triplet model. Add to the gluino axion model two triplet superfields:

$$\hat{\xi}_1 = (\xi_1^{++}, \xi_1^{+}, \xi_1^{0}) : \quad R = 0, \tag{2}$$

$$\hat{\xi}_2 = (\xi_2^{0}, \xi_2^{-}, \xi_2^{--}) : \quad R = +2, \tag{3}$$

then the superpotential (which is required to have $R = +2$) has the following additional terms:

$$\Delta\hat{W} = m_\xi \hat{\xi}_1 \hat{\xi}_2 + f_{ij} \hat{\xi}_1 \hat{L}_i \hat{L}_j + h \hat{\xi}_2 \hat{H}_u \hat{H}_u. \tag{4}$$

Note that the term $\hat{\xi}_1 \hat{H}_d \hat{H}_d$ is forbidden. The resulting scalar potential has the term $|m_\xi \xi_1 + h H_u H_u|^2$, hence the desired trilinear scalar interaction $h m_\xi \xi_1^\dagger H_u H_u + h.c.$ is there to combine with the Yukawa interaction $f_{ij} \xi_1 L_i L_j + h.c.$ to form the well-known dimension-5 effective operator[7] which generates the neutrino masses:

$$(m_\nu)_{ij} = 2 f_{ij} h \frac{\langle H_u \rangle^2}{m_\xi}. \tag{5}$$

If the intermediate scale m_ξ is assumed to be of order the $U(1)_R$ breaking scale, i.e. 10^{11} GeV or so, then m_ν of order 1 eV is obtained if $f_{ij} h$ is of order 10^{-2}.

Consider next the canonical sesaw model. Add to the gluino axion model the singlet superfield $\hat{N}$ with $R = +1$, then the superpotential is supplemented by

$$\Delta\hat{W} = m_N \hat{N} \hat{N} + f_i \hat{L}_i \hat{N} \hat{H}_u \ , \tag{6}$$

which generates the well-known seesaw neutrino mass

$$(m_\nu)_{ij} = f_i f_j \frac{\langle H_u \rangle^2}{m_N} \ . \tag{7}$$

Since both $\hat{N}$ and $\hat{S}$ have the same $U(1)_R$ charge, it is tempting to identify them as one, so that its scalar component has a large vacuum expectation value (VEV) and contains the axion, while its fermionic component is the heavy neutrino singlet of mass m_N. However, the resulting scalar potential will now contain the term $|2m_N\tilde{N}+f_i\tilde{L}_iH_u|^2$, so that the scalar bilinear term $\tilde{L}_iH_u$ (which violates lepton number) has the huge coefficient $2f_im_N\langle\tilde{N}\rangle$ which is clearly unacceptable. To prevent $\hat{N}$ from picking up any VEV, we introduce the discrete symmetry L parity, under which $\hat{L}$, $\hat{e}^c$, and $\hat{N}$ are odd and all other superfields are even, including $\hat{S}$.

In proposing the gluino axion model[2], the composite operator $\mu(\hat{S}) \equiv (\hat{S})^2/M_{Pl}$ with $R=+2$ is used. The couplings of $\mu(S)$ to the supersymmetric particles of the MSSM are required to be invariant under $U(1)_R$. Hence the supersymmetry of the MSSM is broken by $\mu_{eff} = \langle S\rangle^2/M_{Pl}$. In the following we consider an alternative scheme, using the fundamental singlet superfields $\hat{S}_2$, $\hat{S}_1$, and $\hat{S}_0$, with $R=2,1,0$ respectively. We impose the discrete symmetry Z_3 with $\omega^3=1$ on all superfields as follows:

$$1: \qquad \hat{u}^c, \hat{d}^c, \hat{e}^c, \hat{N}, \tag{8}$$

$$\omega: \qquad \hat{Q}, \hat{L}, \hat{S}_1, \hat{S}_0, \hat{\xi}_1 \tag{9}$$

$$\omega^2: \qquad \hat{H}_u, \hat{H}_d, \hat{S}_2, \hat{\xi}_2. \tag{10}$$

We see then that Eqs. (4) and (6) are allowed in addition to Eq. (1) except for the μ term. The superpotential involving $\hat{S}_2$, $\hat{S}_1$, and $\hat{S}_0$ is required to have $R=+2$ also:

$$\hat{W} = m_2\hat{S}_2\hat{S}_0 + f\hat{S}_1\hat{S}_1\hat{S}_0 + h\hat{S}_2\hat{H}_u\hat{H}_d. \tag{11}$$

The resulting scalar potential is

$$V = |m_2S_0 + hH_uH_d|^2 + |2fS_1S_0|^2 + |m_2S_2 + fS_1S_1|^2. \tag{12}$$

Let $v_i \equiv \langle S_i \rangle$, then $V = 0$ has the solution

$$v_0 = 0, \quad v_2 = \frac{-f v_1^2}{m_2}. \tag{13}$$

The problem now is of course the indeterminate value[12] of v_1. To fix v_1 and maintain the above seesaw structure while keeping v_0 zero, we add the following soft terms:

$$V' = -{m'_1}^2 |S_1|^2 - [\lambda m_2 S_2^* S_1^2 + h.c]. \tag{14}$$

The equations of constraint for $V + V'$ to be a minimum are

$$0 = m_2^2 v_2 + (f - \lambda) m_2 v_1^2, \tag{15}$$

$$0 = -{m'_1}^2 + 4f^2 v_0^2 + 2(f - \lambda) m_2 v_2 + 2f^2 v_1^2, \tag{16}$$

$$0 = v_0 (m_2^2 + 4f^2 v_1^2). \tag{17}$$

From Eq. (15), we find

$$v_2 = \frac{(\lambda - f) v_1^2}{m_2}, \tag{18}$$

which indeed preserves the expected seesaw structure. From Eq. (17), we see that $v_0 = 0$ is still a solution, and from Eq. (16), taking into account Eq. (18), we find

$$v_1^2 = \frac{{m'_1}^2}{2\lambda(2f - \lambda)}, \tag{19}$$

where the denominator must be positive for $V + V'$ to be a minimum. The discrete Z_3 symmetry is broken spontaneously by v_1, hence a possible domain wall problem may appear. However, the Majorana fermion singlet $\tilde{S}_1$ may be given a mass m_1 which breaks the Z_3 symmetry softly but explicitly, thus avoiding such a problem.

Note that the scalars S_0 and S_2 remain heavy with mass m_2, but their VEV's are zero or very small[10, 13]. The global $U(1)_R$ symmetry is broken by $\langle S_1 \rangle$ and $\langle S_2 \rangle$, hence the resulting Nambu-Goldstone boson[14] is given by

$$\frac{(v_1)\sqrt{2} Im S_1 + (2v_2)\sqrt{2} Im S_2}{\sqrt{v_1^2 + 4v_2^2}}. \tag{20}$$

In the couplings of S_2 to the superparticles of the MSSM, the axion enters as S_2 is replaced by

$$v_2 e^{2i\varphi} = v_2 e^{2i\langle\varphi\rangle} \exp\left(\frac{ia\sqrt{2}}{v}\right), \tag{21}$$

where $v = \sqrt{v_1^2 + 4v_2^2}$ and the axion a is given by

$$a = (\sqrt{2}v)[\varphi - \langle\varphi\rangle], \tag{22}$$

with $\langle\varphi\rangle = -\theta_{QCD}/6$. Thus the axion decay constant f_a is $\sqrt{2}v \simeq \sqrt{2}v_1$ but M_{SUSY} of the MSSM is v_2. This is analogous to the DFSZ model[3] with M_{SUSY} replaced by M_W.

In this model, the seesaw condition of Eq. (18) implies that $M_{SUSY} \sim f_a^2/M_U$, where M_U is a large effective mass parameter, i.e. $2m_2/h(\lambda - f)$. The allowed range of values for f_a from astrophysics and cosmology[15] is between 10^9 and 10^{12} GeV. Hence M_U is between 10^{15} and 10^{21} GeV. Neutrino masses are given by either Eq. (5) in the Higgs triplet model, or Eq. (7) in the canonical seesaw model. There is no *a priori* connection between f_a and m_ξ or m_N. However, if they are of the same order of magnitude, then m_ν is inversely proportional to f_a as proposed in the models of Ref.[11].

The laboratory detection[16] of axions depends on the $a \to \gamma\gamma$ coupling, which is proportional to[15]

$$\frac{E}{N} - \left(\frac{2}{3}\right)\frac{4 + m_u/m_d + m_u/m_s}{1 + m_u/m_d + m_u/m_s} = \frac{E}{N} - 1.92 \pm 0.08, \tag{23}$$

where N and E are coefficients proportional to the color and electromagnetic anomalies of the axion. For the gluino axion model, $N = 6$ but $E = 0$ without or with neutrino mass from either the canonical seesaw or the Higgs triplet mechanism. This comes from the fact that, except for the gluino, every left-handed fermion has a right-handed partner of the same R.

In conclusion, we have incorporated neutrino masses (through the canonical seesaw or Higgs triplet mechanism) into the gluino axion model, using the superpotentials of Eq. (1)

[without the μ term] and Eq. (11) with either Eq. (6) or Eq. (4). The μ term is replaced by $h\hat{S}_2$, so that

$$\mu_{eff} = h\langle S_2\rangle = hv_2 e^{-i\theta_{QCD}/3}. \tag{24}$$

Assuming that the intermediate scales, i.e. v_1 and m_N or m_ξ are of the same order, we then have

$$M_{SUSY} \sim |\mu_{eff}| \sim \frac{f_a^2}{M_U}, \quad m_\nu \sim f_{eff}^2 \frac{M_W^2}{f_a}, \tag{25}$$

where f_{eff} is a dimensionless coupling.

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References

[1] R. Peccei and H. Quinn, Phys. Rev. Lett. **38**, 1440 (1977); S. Weinberg, Phys. Rev. Lett. **40**, 223 (1978); F. Wilczek, Phys. Rev. Lett. **40**, 279 (1978).

[2] D. A. Demir and E. Ma, hep-ph/0004148.

[3] M. Dine, W. Fischler, and M. Srednicki, Phys. Lett. **B104**, 199 (1981); A. Zhitnitskii, Sov. J. Nucl. Phys. **31**, 260 (1980).

[4] J. E. Kim, Phys. Rev. Lett. **43**, 103 (1979); M. A. Shifman, A. Vainshtein, and V. Zakharov, Nucl. Phys. **B166**, 493 (1980).

[5] C. G. Callan, R. Dashen, and D. J. Gross, Phys. Lett. **63**, 334 (1976); R. Jackiw and C. Rebbi, Phys. Rev. Lett. **37**, 172 (1976).

[6] P. G. Harris *et al.*, Phys. Rev. Lett. **82**, 904 (1999).

[7] S. Weinberg, Phys. Rev. Lett. **43**, 1566 (1979).

[8] E. Ma, Phys. Rev. Lett. **81**, 1171 (1998).

[9] M. Gell-Mann, P. Ramond, and R. Slansky, in *Supergravity*, edited by P. van Nieuwenhuizen and D. Z. Freedman (North-Holland, Amsterdam, 1979), p. 315; T. Yanagida, in *Proceedings of the Workshop on the Unified Theory and the Baryon Number in the Universe*, edited by O. Sawada and A. Sugamoto, KEK Report No. 79-18 (KEK, Tsukuba, Japan, 1979), p. 95; R. N. Mohapatra and G. Senjanovic, Phys. Rev. Lett. **44**, 912 (1980).

[10] E. Ma and U. Sarkar, Phys. Rev. Lett. **80**, 5716 (1998).

[11] R. N. Mohapatra and G. Senjanovic, Z. Phys. **C17**, 53 (1983); P. Langacker, R. D. Peccei, and T. Yanagida, Mod. Phys. Lett. **A1**, 541 (1986); M. Shin, Phys. Rev. Lett. **59**, 2515 (1987); Erratum: **60**, 383 (1988).

[12] E. Ma, Mod. Phys. Lett. **A14**, 1637 (1999).

[13] X. Li and E. Ma, J. Phys. **G23**, 885 (1997).

[14] Y. Nambu, Phys. Rev. Lett. **4**, 380 (1960); J. Goldstone, Nuovo Cimento **19**, 154 (1961).

[15] G. G. Raffelt, Ann. Rev. Nucl. Part. Sci. **49**, 163 (1999).

[16] L. J. Rosenberg and K. A. van Bibber, Phys. Rep. **325**, 1 (2000).