Left-Right Symmetry Just Beyond MSSM, Electric Dipole Moment of the Neutron and HERA Leptoquarks

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Abstract. The supersymmetric left-right (SUSYLR) models solve many of the problems of the minimal supersymmetric standard model (MSSM) such as the R-parity and SUSYCP problems. The first one implies that supersymmetry can provide a naturally stable LSP (lightest supersymmetric particle) which can serve as the cold dark matter of the universe. It is then shown that if $W_R$ mass in the TeV range, the SUSYLR models also provide a natural solution to the strong CP problem without the need for an axion. It is therefore argued that the theory just beyond the MSSM is the SUSYLR model. A crucial prediction of the model is the prediction of the electric dipole moment of the neutron of $d_n \simeq 10^{-25} - 10^{-26}$ ecm arising from one-loop contribution to the strong CP parameter $\bar{\theta}$. Other predictions are a light doubly charged scalar boson and its fermionic superpartner with masses in the few hundred GeV range. Finally, it is pointed out how a simple extension of the model incorporates leptoquarks that can explain the HERA anomaly without giving up R-parity violation

1. Introduction

One of the fundamental new symmetries of nature that has been the subject of intense discussion in particle physics of the past decade is the symmetry

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between bosons and fermions, known as supersymmetry. In addition to the obvious fact that it provides the hope of an unified understanding of the two known forms of matter, the bosons and fermions, it has also provided a mechanism to solve two conceptual problems of the standard model, viz. the possible origin of the weak scale as well as its stability under quantum corrections. The recent developments in strings, which embody supersymmetry in an essential way also have the potential to lead to an ultimate theory of everything. It is therefore naturally a common belief among particle physicists that the next step beyond the standard model is the supersymmetric standard model where the supersymmetry of the model is softly broken by some hidden sector mechanism. The nature and origin of supersymmetry breaking is not relevant for the discussion of what follows. We will therefore make only minimal assumptions about it in this article.

We will first review the basic features of the minimal supersymmetric extension of the standard model (MSSM) and argue that while two of the most critical problems of the standard model i.e. the stability of the Higgs mass and the generation of electroweak symmetry breaking are solved in the MSSM, it creates several new problems which standard model solved in an elegant manner. The new problems include arbitrary amount of baryon and lepton number violation, large electric dipole moment of the neutron among others. The former unfortunately means that MSSM does not provide a natural candidate for cold dark matter of the universe, that is generally considered another major selling point for supersymmetry. We then argue that if the MSSM is assumed to arise as the low energy limit of a supersymmetric left-right model (SUSYLR), then R-parity arises as a natural symmetry keeping baryon and lepton number as automatic symmetries as in the standard model and furthermore it also provides a natural solution to the large electric dipole moment problem of the MSSM (also known as the SUSYCP problem). The lightest supersymmetric particle (LSP) is then a naturally stable particle and can act as the CDM of the universe. More interestingly, we find that if the scale of the righthanded symmetry is in the TeV range, the model also provides a solution to the strong CP problem without the need for an axion, thereby solving a problem of both the standard model as well as the MSSM. A crucial prediction of the theory is that the dipole moment of the neutron arising from the one loop finite contribution to the $\theta$ parameter is of order $10^{-25} - 10^{-26}$ ecm. This is accessible to the next generation of experiments searching for the edm of neutron. This idea therefore has a very good chance to tested in the next decade.
Table 1. The particle content of the supersymmetric standard model. For matter and Higgs fields, we have shown the left-chiral fields only. The right-chiral fields will have a conjugate representation under the gauge group.

| Superfield      | Particles     | Superpartners | gauge transformation |
|-----------------|---------------|---------------|---------------------|
| Quarks $Q$      | $(u, d)$      | $(\bar{u}, \bar{d})$ | $(3, 2, \frac{1}{3})$ |
| Antiquarks $U^c$| $\bar{u}^c$  | $\bar{u}^c$ | $(3^*, 1, -\frac{4}{3})$ |
| Antiquarks $D^c$| $\bar{d}^c$  | $\bar{d}^c$ | $(3^*, 1, \frac{1}{3})$ |
| Leptons $L$     | $(\nu, e)$   | $(\bar{\nu}, \bar{e})$ | $(1, 2 - 1)$ |
| Antileptons $E^c$| $\bar{e}^c$ | $\bar{e}^c$ | $(1, 1, 2)$ |
| Higgs Boson $H_u$ | $(H_u^+, H_u^0)$ | $(\bar{H}_u^+, \bar{H}_u^0)$ | $(1, 2, +1)$ |
| Higgs Boson $H_d$ | $(H_d^0, H_d^-)$ | $(\bar{H}_d^0, \bar{H}_d^-)$ | $(1, 2, -1)$ |
| Color Gauge Fields | $G_a$ | $\bar{G}_a$ | $(8, 1, 0)$ |
| Weak Gauge Fields | $W^\pm, Z$ | $\bar{W}^\pm, \bar{Z}$ | |
| Photon          | $\gamma$     | $\bar{\gamma}$ | |

2. The MSSM and its problems

The MSSM is the minimal supersymmetric extension of the standard model and is based on the same gauge group as the standard model i.e. $SU(3)_c \times SU(2)_L \times U(1)_Y$. In Table 1, we give the particle content of the model.

The first point to note is that while the gauge interaction of the standard model fermions remains unchanged in this supersymmetrized version, the weak interactions of the squarks and the sleptons are very different from their fermionic partners due to supersymmetry breaking. This has the phenomenological implication that the gaugino-fermion-sfermion interaction changes generation leading to potentially large flavor changing neutral current (FCNC) effects such as $K^0-\bar{K}^0$ mixing, $\mu \rightarrow e\gamma$ decay etc unless the sfermion masses of different generations are chosen to be very close in mass. This is the FCNC problem of the MSSM and we do not dwell on this here since the SUSY-LR model does not throw light on this.

To discuss the other problems of MSSM, let us discuss the superpotential of the model. It consists of two parts:

$$W = W_1 + W_2,$$

where

$$W_1 = h^{ij}_t E_i^c L_j H_d + h^{ij}_d Q_i D_j^c H_d + h^{ij}_e Q_i U_j^c H_u + \mu H_u H_d$$

$$W_2 = \lambda_{ijk} L_i L_j E_k^c + \lambda'_{ijk} L_i Q_j D_k^c + \lambda''_{ijk} U_i^c D_j^c D_k^c,$$
\(i, j, k\) being generation indices. Note that while the terms in \(W_1\) conserve baryon and lepton number those in \(W_2\) do not and yet they are perfectly allowed by the gauge invariance of the model. Thus baryon and lepton number are no more natural symmetries of the model as they were for the standard model. The latter are known as the \(R\)-parity breaking terms where \(R\)-parity is defined as \(R = (-1)^{3(B-L)+2S}\) where \(S\) is the spin of the particle. These \(R\)-parity violating couplings are severely constrained by present experiments. For instance, present experimental limits on the lifetime of the proton imply that the product \(\lambda \lambda'' \leq 10^{-24}\) which is much too severe a fine tuning. Furthermore the LSP in this model is too shortlived to play the role of CDM of the universe.

In order to prevent the appearance of these terms, one generally assumes the existence of some global symmetry (which could be discrete or continuous). First of all this is an additional assumption which diminishes the appeal of MSSM somewhat; but perhaps more serious is the lore that nonperturbative quantum gravity tends to break all global symmetries of nature making the appearance of Planck suppressed \(\Delta B \neq 0\) and \(\Delta L \neq 0\) in the low energy Lagrangian quite plausible. Even though in general the strengths of these terms are small due to the Planck scale suppression, they are not small enough to let the LSP live longer than the age of the universe. Thus it may be preferable to search for more natural ways to eliminate the \(R\)-parity violating terms from the supersymmetric models. It was realized in mid-80’s that such is the case in the supersymmetric version of the left-right model that implements the see-saw mechanism. We demonstrate this point in the next section.

A second problem with the MSSM lies in its predictions for the CP violating effects being too large. To see this note that in the simple versions of MSSM, there are at least three phases including the usual KM-phase \(\delta_{KM}\) of the standard model. The two of the extra phases reside in the soft breaking parameter \(A\) and the Higgs mixing mass \(\mu\). A priori, all phases in a theory are arbitrary. In the standard model even though the KM phase can be left arbitrary, its physical effects are always suppressed by the product of various quark mixing angles leading to natural suppression of all CP violating effects. In the MSSM however, such suppressions do not arise in all CP violating effects. For instance, if we calculate the electric dipole moment of the neutron, then we typically get an expression of the form:

\[
d_n^e \simeq \frac{e}{16\pi^2} \frac{m_d}{M_q^2} \text{Arg}(m_{\tilde{g}}[A - \mu \tan \beta])
\]

A simple evaluation of the above down quark electric dipole moment leads to the conclusion that unless either (i) the squark masses are of order 3 TeV or (ii) \(\text{Arg}(m_{\tilde{g}}A)\) and \(\text{Arg}(m_{\tilde{g}}\mu)\) are less than \(10^{-3}\) if squark masses \(M_q \simeq 100\) GeV, the edm of the neutron will come out to be three orders
of magnitude higher than the present experimental upper bound. In either case, we have a fine tuning problem for the theory, the very problem supersymmetry was supposed to solve. In the first case one has to fine tune to get the Higgs mass of order $m_W$ and in the second case, the new phases of the model (unlike the CP phase of the standard model) has to be tuned down by three orders of magnitude from its natural value.

The third and final problem of MSSM that we discuss here is the strong CP problem. It is by now quite well-known that the Quantum Chromodynamics theory of strong interactions which has been so successful in its description of low and high energy hadronic phenomena implies an arbitrary amount of CP-violation in the $\Delta S = 0$ processes. This results from the periodic vacuum structure of QCD and is given by the following effective Lagrangian:

$$L_{\theta} = \theta \frac{g^2}{32\pi^2} G_{\mu\nu} \tilde{G}^{\mu\nu}$$  \hspace{1cm} (5)

This has the disastrous implication that it leads to an electric dipole moment of the neutron given by $d_n \approx 10^{-16} \theta$ ecm. The present upper limits on the $d_n \leq 10^{-25}$ ecm then imply that $\theta \leq 10^{-9}$ which requires a severe fine tuning of the parameters and is therefore not acceptable in a final theory. The inclusion of the electroweak sector of the theory leads to a slight modification of the above results: the actual observable $\bar{\theta}$ is called $\bar{\theta} \equiv \theta + \text{ArgDet} M_q$. The above upper bound actually applies to $\bar{\theta}$.

What goes wrong if we set $\bar{\theta} = 0$ by hand in a given theory? The electroweak sector of the theory then generates an infinite contribution to $\bar{\theta}$ in loop levels requiring that we set the $\bar{\theta}$ to a small value in every order of perturbation theory. For instance, in the standard model, one finds a nonzero and infinite $\bar{\theta}$ at the sixth loop level. Coming to the MSSM, since one now has colored gluinos, the effective parameter describing strong CP violation is now given by:

$$\bar{\theta} = \theta + \text{ArgDet} M_q - 3\text{Arg} M_{\tilde{G}}$$  \hspace{1cm} (6)

We see that the same gluino phase that was responsible for large edm of neutron at the one loop level now contributes also to the $\bar{\theta}$ at the tree level. If we set the gluino phase to zero at the tree level by hand, we will get an infinite value for the phase in the MSSM at the two loop level (see Figure 1). Thus one could argue that in the MSSM, the strong CP problem is worse than in the standard model.

The most popular solution to the strong CP problem is the Peccei-Quinn solution which requires the complete gauge theory of electroweak and strong interactions to respect a global chiral $U(1)$ symmetry. This symmetry must however be spontaneously broken in the process of giving mass to the $W$-boson and fermions leading to a pseudo-Goldstone boson
in the particle spectrum known in the field as the axion. There are three potential problems with this otherwise beautiful proposal: (i) the axion has not been experimentally discovered as yet and the window is closing on it; and (ii) if non-perturbative gravitational effects induced by black holes and wormholes are important in particle physics as is believed by some, then the axion solution would require fine tuning of the gravitationally induced couplings by some 50 orders of magnitude. This will make the axion theory quite contrived. Finally, it is not easy to get the correct scale for the axion in superstring models which in the opinion of many people is the ultimate theory of all matter and forces.

A second class of solutions that does not lead to any near massless boson is to require the theory to be invariant under discrete symmetries. The most physically motivated of such theories are the ones based on the left-right symmetric theories of weak interactions. These theories are based on the gauge group $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$ with quarks and leptons assigned in a left-right symmetric manner. Such models are also completely quark-lepton symmetric. To see how parity symmetry of the Lagrangian helps to solve the strong CP problem, let us note that invariance under parity sets $\theta = 0$ because $G \tilde{G}$ is odd under parity. Additionally, constraints of left-right symmetry imply that the Yukawa couplings of quarks are hermitean. If furthermore the vacuum expectation values (VEVs) of the Higgs fields responsible are shown to be real, then this would automatically lead to $\bar{\theta} = 0$ at the tree level. If the one loop contributions also preserve the hermiticity of the quark mass matrices, then we have a solution to the strong CP problem. In the nonsupersymmetric left-right models with nontrivial CP violation, it is well known that in general VEVs of the Higgs field are not real. This, in the past led to suggestions that either new discrete symmetries be invoked together with left-right symmetry or new vectorlike fermions be added to the theory. Such theories also do not suffer from the Planck scale implied fine tunings. It always remained a challenge to solve the strong CP problem using only left-right symmetry since often new additional symmetries invoked are not motivated from any other consideration. Recently it has been shown that left-right symmetry in combination with supersymmetry solves the strong CP problem without the axion.

It is the goal of this talk to discuss how the dark matter, SUSY CP as well as the strong CP problems are all solved without any extra assumptions and introducing any extra fermions if the MSSM arises as a low energy limit of the supersymmetric left-right (SUSYLR) model with the Higgs structure required to obtain the see-saw mechanism for the neutrino masses. Furthermore solution to the strong CP problem requires that the scale of $SU(2)_R$ symmetry breaking must be in the TeV range. A crucial prediction of the model is a dipole moment of the neutron $10^{-25} - 10^{-26}$ ecm, which
in the range accessible to the next generation of the proposed experiments.

3. Supersymmetric Left-Right model

The left-right symmetric theories are based on gauge group \( SU(2)_L \times SU(2)_R \times U(1)_{B-L} \) with quarks and leptons transforming as doublets under \( SU(2)_L, R \). These models were originally proposed\[12\] in order to understand the origin of parity violation as a consequence of spontaneous symmetry breaking rather than as an intrinsic part of the gauge interactions. It was then used to develop models of CP violation as well as to discuss the strong CP problem\[10\]. The version of the model we will be interested in was proposed by G. Senjanović and this author \[14\] in order to implement the see-saw mechanism to understand the small neutrino masses. This see-saw requirement fixes the Higgs structure of the model and this is the version which seems to have all the desirable properties that we are interested in here.

In Table 2, we denote the quark, lepton and Higgs superfields in the theory along with their transformation properties under the gauge group. Note that we have chosen two bidoublet fields to obtain realistic quark masses and mixings (one bidoublet implies a Kobayashi-Maskawa matrix proportional to unity, because supersymmetry forbids \( \tilde{\Phi} \) in the superpotential).

The superpotential for this theory is given by (we have suppressed the generation index):

\[
W = Y_q^{(i)} Q^T \tau_2 \Phi_1 \tau_2 Q^c + Y_l^{(i)} L^T \tau_2 \Phi_1 \tau_2 L^c \\
+ i(f L^T \tau_2 \Delta L + f^c L^c \tau_2 \Delta^c L^c)
\]
\[ + \mu_\Delta \text{Tr}(\Delta \bar{\Delta}) + \mu_\Delta^\epsilon \text{Tr}(\Delta^\epsilon \bar{\Delta}^\epsilon) + \mu_{ij} \text{Tr}(\tau_2 \Phi_i^T \tau_2 \Phi_j) \]
\[ + \text{W}_{NR} \quad (7) \]

where \( W_{NR} \) denotes non-renormalizable terms arising from higher scale physics such as grand unified theories or Planck scale effects. At this stage all couplings \( Y_{q_1}, \mu_{ij}, \mu_{\Delta}, \mu_{\Delta^\epsilon}, f_c, f \) are complex with \( \mu_{ij} \), \( f \) and \( f_c \) being symmetric matrices.

The part of the supersymmetric action that arises from this is given by
\[ S_W = \int d^4x \int d^2\theta W + \int d^4x \int d^2\bar{\theta} W^\dagger. \quad (8) \]

The terms that break supersymmetry softly to make the theory realistic can be written as
\[ \mathcal{L}_{\text{soft}} = \int d^4 \theta \sum_i m_i^2 \phi_i^\dagger \phi_i + \int d^2 \theta \theta^2 \sum_i A_i W_i + \int d^2 \bar{\theta} \bar{\theta}^2 \sum_i A_i^* W_i^\dagger \]
\[ + \int d^2 \theta \theta^2 \sum_p m_p \bar{W}_p W_p + \int d^2 \bar{\theta} \bar{\theta}^2 \sum_p m_p^* \bar{W}_p^* W_p^*. \quad (9) \]

In Eq. 9, \( \bar{W}_p \) denotes the gauge-covariant chiral superfield that contains the \( F_{\mu\nu} \)-type terms with the subscript going over the gauge groups of the theory including SU(3)\(_c\). \( W_i \) denotes the various terms in the superpotential, with all superfields replaced by their scalar components and with coupling matrices which are not identical to those in \( W \). Eq. 9 gives the most general set of soft breaking terms for this model.

It is clear from the above equations that this model has no baryon or lepton number violating terms and if the ground state of the theory does not break \( B \) or \( L \), then the LSP in this model is stable and becomes the dark matter particle. We will address the issue of the ground state of the theory in a subsequent section since the requirement of R-parity conservation in the ground state requires that the model have in it nonrenormalizable operators induced by higher scale physics such as Planck scale effects.

We noted earlier that left-right symmetry implies that the first term in \( \bar{\theta} \) is zero. Let us now see how supersymmetric left-right symmetry also requires the second term in this equation to vanish naturally. Under left-right transformation, the fields and the supersymmetric variable \( \theta \) transform as follows:
\[ Q \leftrightarrow Q^{c*} \]
\[ L \leftrightarrow L^{c*} \]
\[ \Phi_i \leftrightarrow \Phi_i^\dagger \]

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\[ \Delta \leftrightarrow \Delta^c \]
\[ \Delta \leftrightarrow \Delta^c \]
\[ \theta \leftrightarrow \bar{\theta} \]
\[ \tilde{W}_{SU(2)_{L}} \leftrightarrow \tilde{W}^{*}_{SU(2)_{R}} \]
\[ \tilde{W}_{B-L,SU(3)_{C}} \leftrightarrow \tilde{W}^{*}_{B-L,SU(3)_{C}} \]

(10)

Note that this corresponds to the usual definition \( Q_{L} \leftrightarrow Q_{R} \), etc. With this definition of L-R symmetry, it is easy to check that

\[ Y_{q,i}^{(i)} = Y_{q,i}^{(i)\dagger} \]
\[ \mu_{ij} = \mu_{ij}^{*} \]
\[ \mu_{\Delta} = \mu_{\Delta}^{*} \]
\[ f = f_{c}^{*} \]
\[ m_{\lambda_{SU(2)_{L}}} = m_{\lambda_{SU(2)_{R}}}^{*} \]
\[ m_{\lambda_{B-L,SU(3)_{C}}} = m_{\lambda_{B-L,SU(3)_{C}}}^{*} \]
\[ A_{i} = A_{i}^{\dagger} \]

(11)

The first point to note is that the gluino mass is automatically real in this model; as a result, the last term in the equation for \( \bar{\theta} \) above is naturally zero. We now therefore have to investigate only the quark mass matrices in order to guarantee that \( \bar{\theta} \) vanishes at the tree level. For this purpose, we note that the Yukawa matrices are hermitean and the mass terms involving Higgs bidoublets in the superpotential are real. If we can show that the vacuum expectation values of the bi-doublets are real, then the tree level value of \( \bar{\theta} \) will be naturally zero. For this purpose note that the second observation above implies that the Higgs potential involving the bidoublet and the triplet Higgs fields is completely CP conserving (i.e. it has no complex parameters in it). This is encouraging since a completely CP-conserving potential always has a domain of parameters where the ground state is CP conserving. The story however is not so simple since the soft breaking terms involving the sneutrinos (e.g. \( \tilde{\nu} \phi \tilde{\nu}^{c} \)) have a complex coupling and if the ground state has a nonzero sneutrino vev, the ground state will break CP. A careful study of the ground state is therefore essential for this purpose.

4. **Reality of the bidoublet vev’s, tree level \( \bar{\theta} \) and limit on the \( W_{R} \) scale.**

To see whether the bidoublet vevs are real, one has to investigate the ground state of the theory and in particular we have to see whether the sneutrinos
have vev. This was investigated by R. Kuchimanchi and this author in Ref. [18]. It was noted there that the ground state of the SUSYLR model does violate common intuition i.e. if only renormalizable terms in the superpotential are kept in the minimal SUSYLR model, the ground state does not break any of the gauge symmetries nor parity symmetry. This is easily seen in the supersymmetric limit since in this case the mass-squares of the $\Delta$ Higgses are positive as are those for the $\Phi$ fields and since the D-term contributions to the potential are always positive, the above result follows. Once supersymmetry breaks, in principle some of the above mass-squares can be negative and the argument is not so simple. It however turns out that if a particular Higgs mass-square is negative, there are particular directions in the Higgs space along which the D-term vanishes and in those directions the negative mass-square will make the potential unbounded from below. This means that no mass square should be allowed to be negative. This of course means that the symmetry cannot break. The details of this argument are given in ref. [18].

This result can however be avoided by adding a parity odd singlet $\sigma$ to the model in which case one can write a superpotential of the form

$$\begin{align*}
W(\sigma, \Delta...) &= \mu(\Delta\Delta + \Delta^c\Delta^c) + \lambda\sigma(\Delta\Delta - \Delta^c\Delta^c) + m_\sigma\sigma^2
\end{align*}$$

which after supersymmetry breaking soft terms gives rise to parity violating and SU$(2)_R$ violating minimum. The problem now is that[18] is that the ground state of this slightly extended model is also quite unusual in that it breaks electric charge unless R-parity is spontaneously broken by the $<\tilde{\nu}^c> \neq 0$. But this is precisely what we wanted to avoid in order to get real bidoublet vev’s and solve the strong CP problem. How is this to be achieved? It was pointed out in [16] that if one adds nonrenormalizable Planck scale induced terms to the superpotential, one can indeed have a parity violating vacuum that conserves R-parity. Since this is such an important point, we review this proof in an appendix. The conclusion of this part then is that if one includes non-renormalizable Planck scale induced terms to the minimal SUSYLR model, the ground state corresponds to having zero sneutrino vev’s and therefore all couplings that zeroth order in $M^{\frac{3}{2}}$ in the Higgs potential are real leading to real bidoublet vevs to this order. Note however that one must include all allowed nonrenormalizable terms to the model and one such term has the form $\phi\phi\Delta^c\Delta^c/M_{Pl}$ which can come with a complex coupling. This will then induce a tree level phase in the bidoublet vevs of order $v^3_{R}/M_{Pl}^2$ [19]. For its effect on $\bar{\theta}$ to be less than $10^{-10}$, one must have $v_R \leq 10^5$ GeV or so.

To summarize this sub-section, the requirement of tree level value of theta to be naturally small implies that there must be non-renormalizable terms in the superpotential and the scale of the right handed $W_R$ must be in the few TeV range.
5. One-loop contributions to $\bar{\theta}$ and the electric dipole moment of the neutron

Having shown that the tree level contributions to the $\bar{\theta}$ are small, we now proceed to discuss the magnitude of the one loop contributions to the quark and gluino masses. There are two classes of diagrams that have to be considered for the quark masses: one where the gluino intermediate state is present and another where the Wino states contributes. Let us first discuss the gluino effect given in the Fig.2.

5.1. Gluino contribution

We work for simplicity in the case in which the vevs of $\Phi_a$ are in the form

$\text{Diag } <\Phi_1> = (\kappa_u,0)$ and $\text{Diag } <\Phi_2> = (0,\kappa_d)$. Note that due to left-right symmetry the gluino mass is real. Any complexity of the determinant of the quark mass matrices must therefore come from the supersymmetry breaking contributions to the squark masses. The general form of the squark mass matrix for the down sector can be written as

$$M^2_{\tilde{d} \tilde{d}} = \begin{pmatrix} m^2_L + M^2_d + c_u h_u^2 & X \\ X^\dagger & m^2_R + M^2_d + c'_u h_u^2 \end{pmatrix}$$

(13)

where $X = (A - \mu \tan \beta)(M_d + a_u M_d h^2_u + h^2_u M_d a'_u)$. We have ignored small contributions proportional to the down quark couplings. Note that in the limit of exact left-right symmetry, $c_u = c'_u$, $a_u = a'_u$. However, after parity breaking effects are included, we have $c_u - c'_u \simeq \frac{\alpha_s}{16\pi} \ln \frac{m^2_u}{m^2_W}$, etc. In order to calculate the one loop contribution to the quark edm, we focus on the down quark and treat the $c$ terms, the $M^2_d$ and the $X$-terms as perturbations. To zeroth order in this perturbation, then we have a diagram which is given in figure 2 but without the solid circles which represent the insertion of $c_u$ and $c'_u$ terms. Also let us work in a basis in which the up quark mass matrix is diagonal. It is then easy to see that, the loop is proportional to $a_u M_d h^2_u + a'_u h^2_u M_d$. Since $M_d$ is hermitean and $M^2_u$ is diagonal, the 11 element of this matrix is real and therefore it makes no contribution to the edm of the down quark. In order to see at what level we will get a contribution to the quark edm we keep inserting the various perturbation terms proportional to $c_u$ and $c'_u$ until we get a complex 11 entry for the down quark mass matrix. It turns out that the dominant contribution to the mass matrix is given by

$$M^d_{\text{loop}} \simeq \text{Im}(V^*_{tq} V_{tb} V_{cd} V^*_{cd}) \frac{\alpha_w}{4\pi} \frac{3\alpha_W}{2\pi} h^2_u h^2_t m_b Z_{\text{susy}} \ln \frac{M^2_{W_R}}{M^2_{W_L}}$$

(14)

where $Z_{\text{susy}} = m_c (A - \mu \tan \beta)$. The fact that this CP-violating contribution would be proportional to the product of the CKM mixing angles was noted.
in [16] from simple analytic arguments. In ref.[20], an enhancement of the result by a factor \( m_b/m_d \) was noted. Thus the contribution of this to the \( \bar{\theta} \) is:

\[
\bar{\theta} \simeq \text{Im}(V^{*}_t V^*_b V^*_c V_{cd}) \frac{\alpha_s}{4\pi} \frac{3\alpha_W}{2\pi} h_c^2 m_b \frac{m_b}{m_d} Z_{\text{susy}} \ln \frac{M^2_{WR}}{M^2_{WL}}
\] (15)

where we have set \( h_t \simeq 1 \) as is suggested by the recent top quark mass measurements. Putting in all the numbers, one finds that \( \bar{\theta} \simeq 10^{-9} \). Barring unforeseen cancellation of parameters therefore, this will lead to a value for the electric dipole moment of the neutron of order \( 10^{-25} \) to \( 10^{-26} \) ecm for \( W_R \) masses in the TeV range. This number is quite close to the present upper limit. At present there are plans to improve the search for the edm of neutron by several orders of magnitude[21] by using ultra cold neutron in superfluid He\(^4\). These experiments will therefore provide a decisive test of this class of supersymmetric left-right models.

5.2. One loop wino contributions and C invariance

Next let us consider the wino contributions to the quark masses. A typical graph is shown in Fig.3. We see from this that it involves the majorana mass for the SU(2)\(_L\) wino. There is a similar graph involving the SU(2)\(_R\) wino. If parity was an exact symmetry and the full mass of the left and the right winos were the same, these two two graphs would have added up to give a real quark mass at the one loop. However since parity breaking leads to a non-cancellation between these two diagrams, we will get a contribution to quark masses of order \( \simeq \frac{\alpha_w}{4\pi} M_{\lambda_w} \). In theories where the wino and zino masses vanish, this contribution is suppressed. However in general, this will be present. Again it may turn out that in specific theories, this mass may be real (as in some gauge mediated supersymmetry breaking models) in which case again, this keeps the tree level hermiticity of the mass matrix and the contribution to \( \theta \) vanishes. In general however, the dynamical content of the theory may be such that neither happens. In this case, it is necessary to impose in addition a charge conjugation (C) symmetry on the Lagrangian which makes these contributions vanish[19]. It actually turns out that the C-symmetry must be softly broken to keep the CKM phase in weak current from vanishing.

5.3. One loop phase in the gluino masses

A rough order of magnitude of the CP violating phase in the gluino mass can be estimated as follows: since the \( A_{u,d} \) are hermitean and proportional to the Yukawa couplings \( h_{u,d} \) at some scale above the \( M_R \) scale, let us go to a basis where \( h_d \) and \( A_d \) are diagonalized. Then we find that, at the scale of proportionality, if any one of the off-diagonal elements of \( h_u \) and
(hence $A_w$) are set to zero, the theory becomes completely CP conserving and cannot generate a CP violating phase at any scale below $M_R$. It is then clear that the one loop graph that generates a phase in the gluino mass can lead to the gluino phase $\delta_{\tilde{g}}$ which is at most

$$\delta_{\tilde{g}} \simeq \frac{V_{ub}V_{cb}V_{cd}V_{du}}{8\pi^2}Z_{\text{susy}} \frac{m_\nu^2}{M^2} \ln \frac{M^2}{M^2}$$

leading to $\delta_{\tilde{g}} \leq 10^{-11}$ which is negligible.

Thus the conclusion of this section is that the finite values of $\bar{\theta}$ induced at the one loop level are within the present experimental limits so that SUSYLR model indeed provides a satisfactory solution to the strong CP problem. The SUSYCP problem is then of course automatically solved.

6. Phenomenological Implications: light doubly charged Higgs bosons

The low $W_R$ left-right models are rich in phenomenology as is well known. Apart from the obvious searches for the new particles of the model such as the $W_R$, $Z_{LR}$ and the heavy right-handed neutrino, there are a large number of new processes induced by the non-standard model particles present. A very typical example of such process is the neutrinoless double beta decay induced in these models by the exchange of right-handed neutrinos $[22]$, doubly charged Higgs bosons $[23]$ etc. In fact, the recent stringent lower limits on the lifetime for neutrinoless double beta decay in enriched $^{76}\text{Ge}$ by the Heidelberg-Moscow collaboration $[24]$ implies that $M_{W_R} \geq 1.1$ TeV and $M_{\nu_R} \geq 1.1$ TeV when combined with theoretical limits from vacuum stability. There are also comparable stringent limits on the $W_R$ mass of the order of a TeV from the $K^0 - \bar{K}^0$ mass difference $[25]$. The bidoublet Higgs induced flavor changing neutral currents also imply that the all but the lightest standard model Higgs boson must be in the 5 to 10 TeV range $[26]$. There are also strong limits on the charged Higgs masses $[27]$ and the $W_L - W_R$ mixing angle $[28]$ from the $b \rightarrow s\gamma$ measurements. Finally there are recent collider limits on the mass of the $W_R$ and $Z_{LR}$ $[29]$ from the D0 and CDF experiments $[30]$. The detailed discussion of these topics is beyond the scope of this talk.

Another very striking prediction of these models is the existence of light (100 GeV to TeV range) doubly charged Higgs bosons. These lead to many interesting phenomenological implications such as muonium-antimuonium transition $[31]$ on which there now exist the most stringent upper limit of $\leq 3 \times G_F 10^{-3}$ from the PSI $[32]$. This implies that the combination of couplings and masses of the doubly charged bosons is restricted as follows:

$$\frac{f_{\mu\mu}}{\Delta} \leq 2 \times 10^{-7} \text{GeV}^{-2}.$$  

This particle also contributes to the $g-2$ of
the muon. The present measurements of $g-2$ imply that the additional contribution from the doubly charged Higgs boson must be bounded above by $8.5 \times 10^{-9}$ implying that $\frac{f_{\mu\mu}}{M_\Delta^2} \leq 1.4 \times 10^{-4} GeV^{-2}$. The present BNLE821 experiment aims to push the accuracy to the level of $4 \times 10^{-10}$ level making the above constraint more stringent by a factor of 21. They also have many implications for collider phenomenology.

There have also been study of the new supersymmetric particles of these models, specially in regard to their manifestation in colliders[34]. The literature in the subject is vast and we do not discuss it any further. Instead, I focus on the recent work on trying to explain the HERA anomaly in terms of the leptoquarks that arise naturally in the context of the SUSYLR models in the next section.

7. Leptoquarks and HERA anomaly in the SUSYLR model

If the high $Q^2$ anomaly observed recently in the $e^+p$ scattering by the H1[35] and the ZEUS[36] collaborations is confirmed by future data, it will be an extremely interesting signal of new physics beyond the standard model. A very plausible and widely discussed interpretation of this anomaly appears to be in terms of new scalar particles capable of coupling to $e^+u$ or $e^+d$ of scalar leptoquarks[37] with mass around 200 GeV. Alternative interpretations based on contact interactions[38] or a second $Z'$[39] have been proposed; but attempts to construct models that lead to the desired properties seem to run into theoretical problems.

The leptoquark must be a spin zero color triplet particle with electric charge $5/3$ or $2/3$. The latter electric charge assignment makes it possible to give a plausible interpretation of the leptoquark as being the superpartner of the up-like quark[40] of the minimal supersymmetric standard model (MSSM) provided one includes the R-parity violating couplings of the type $\lambda' Q LD^c$ to the MSSM. There are however very stringent upper limits on several R-parity violating couplings: for instance, if in addition to the $\lambda'$ term described above, one adds the allowed $\lambda'' u^c d^c d^c$ terms to the superpotential, then it leads to catastrophic proton decay unless $\lambda' \lambda'' \leq 10^{-24}$[2]. There are also stringent limits on $\lambda'_{111} \leq 10^{-4}$[11] from neutrinoless double beta decay, which forces the lepto-quark to be $\tilde{c}$ or $\tilde{t}$ rather than the obvious choice $\tilde{u}$. Furthermore, within such a framework, the lightest supersymmetric particle (LSP) is no more stable and therefore, there is no cold dark matter (CDM) candidate in such theories. Since supersymmetric left-right lead to automatic R-parity conservation we have recently used them as a typical framework for studying the consequences of leptoquarks without giving up R-parity conservation[42]. We augment the SUSYLR model by including the leptoquark fields. Demanding that the
lepto-quarks couple to $e^+d$ or $e^+u$ leads to the conclusion that they must belong to the multiplet $(2, 2, 4/3, 3^*)$ (denoted $\Sigma_{QL^e}$). Anomaly cancellation requires that there be a conjugate state $(2, 2, -4/3, 3^*)$ (denoted by $\Sigma_{Q^cL}$). Each of these multiplets have four scalar fields and four fermion fields which will be denoted in what follows by the obvious subscript corresponding to their couplings. We denote the four scalar leptoquarks as $\Sigma_{ue^c}, \Sigma_{u\nu^c}, \Sigma_{de^c}$ and $\Sigma_{d\nu^c}$ and their fermionic partners (to be denoted by a tilde on the corresponding scalar field). Writing down the most general superpotential, one can easily convince oneself that the resulting theory maintains the property of automatic R-parity conservation.

Before discussing the application of this model to explain the HERA anomaly, let us first discuss the mass spectrum of the model. The superpotential for this model will have a direct mass term of the form $M_0\Sigma\overline{\Sigma}$ which will imply that the fermionic fields in the leptoquark multiplet will have a masses $M_0$ prior to symmetry breaking. There may be other contributions to these masses from radiative corrections which will split the degeneracy implied by the above mass term.

As far as the scalar leptoquark states are concerned, their masses will receive several contributions: first a direct common contribution from the $M_0$ term given above. After symmetry breaking, the D-terms of the various gauge groups will contribute. There is also soft SUSY breaking contribution along with the radiative correction. Assuming that the $SU(2)_R$ symmetry is broken by the vev’s $<\Delta^c>=v_1$, $<\overline{\Delta}^c>=v_2$ and the $SU(2)_L$ symmetry is broken by the two $\phi$ vev’s as $\text{diag}<\phi_u>=(0, v\sin\beta/\sqrt{2})$ and $\text{diag}<\phi_d>=(v\cos\beta/\sqrt{2}, 0)$, we can write the masses for the various scalar leptoquarks as follows:

$$M_{\Sigma_{ue^c}}^2 = M_0^2 + (I_R^a g_{2R}^2 - B - L) \frac{(v_1^2 - v_2^2)}{v_2} + \frac{g_{2R}^2 v^2}{4} \cos 2\beta + \frac{\Delta_4^m}{2} + \text{Radiative correction}$$

The values of the $I_R^a$ and $B - L$ are given in Table I:

The radiative correction can be positive or negative. From table I and Eq.(1), we see easily that if $v_1 < v_2$, then the lightest leptoquark state is $\Sigma_{ue^c}$ since $g_{2R}^2 > 2g_{B-L}^2$ in the left-right models. Furthermore, interestingly enough for this choice of vev’s, assuming the combined $\Delta_4^m + \text{Radiative correction}$ to be smaller than the the D-term, the leptoquarkino states are heavier than the lightest leptoquark state. We assume their masses to be in the range of 300 to 400 GeV. In a subsequent section we will obtain a lower bound on the leptoquarkino mass from the present collider data.

Turning to the couplings of the leptoquarks $\Sigma$ and $\overline{\Sigma}$ to quarks and
leptons, it is given by the superpotential:

\[ W_{lq} = \lambda_{ij} (\Sigma Q_i^c L_j + \bar{\Sigma} Q_j^c \bar{L}_i) \]  

(18)

(where \( i, j \) are generation indices). Let us assume for simplicity that \( \lambda_{ij} \) are diagonal in the mass eigenstate basis for the quarks and leptons. Explanation of the HERA high \( Q^2 \) anomaly seems to require \( \lambda_{11} \approx 0.05 \) which we will assume from now on and mass of the scalar field in \( \Sigma \) (assumed to be \( \Sigma_{\nu_e} \) from the above mass arguments) around 200 GeV. As far as the other couplings go, they will be strongly constrained by the present experimental upper limits on the low energy processes such as \( \mu \rightarrow e\gamma, \tau \rightarrow e\gamma \) etc. which will arise at the one loop level from the exchange of \( \Sigma \) and \( \bar{\Sigma} \). There are also tree level diagrams which can lead to rare processes such as \( K \rightarrow \pi e^-\mu^+ \). These processes imply an upper limit of \( \lambda_{22} \leq 2 \times 10^{-3} \). The most stringent upper limit on \( \lambda_{22} \) comes from the upper limit on the process \( K^0_L \rightarrow \mu^+e^- \) and yields \( \lambda_{22} \leq 2 \times 10^{-4.5} \). Turning to \( \lambda_{33} \), the most stringent limits arise from the present upper limit on the branching ratio for the process \( \tau \rightarrow e\gamma \) which the 1996 Particle data tables give as \( \leq 1.1 \times 10^{-4} \). This implies a weak upper limit on \( \lambda_{33} \leq 0.2 \). Thus the third generation leptoquark coupling \( \lambda_{33} \) could in principle be comparable to \( \lambda_{11} \). If the efficiency for the detection of \( \tau \) leptons at HERA were comparable to the detection efficiency for electrons, that could also severely limit \( \lambda_{33} \). Finally we also note that such a value for \( \lambda_{11} \) is also consistent with the present data on the parity violation in atomic physics. Thus it appears that our leptoquark couplings are consistent with all known low energy data.

An important consequence of the above coupling is that the branching ratio for the leptoquark coupling to the electron mode is only 50%. As a result the recent CDF/D0 bound for the leptoquark mass is around 195 GeV[43] which is lower than mass required for explaining the HERA anomaly.

The existence and properties of the leptoquarkino provides a new and unique signature for our model as compared to all other proposals to ex-

| states | \( I^a_{3R} \) | \( B - L \) | \( I_{3L} \) |
|--------|----------------|-------------|------------|
| \( u^e e \) | \( \frac{1}{2} \) | \( \frac{1}{2} \) | \( \frac{1}{2} \) |
| \( u^\nu \) | \( \frac{1}{2} \) | \( \frac{1}{2} \) | \( -\frac{1}{2} \) |
| \( u^e e \) | \( -\frac{1}{2} \) | \( -\frac{1}{2} \) | \( -\frac{1}{2} \) |
| \( d^e e \) | \( -\frac{1}{2} \) | \( -\frac{3}{2} \) | \( \frac{1}{2} \) |
| \( d^\nu \) | \( -\frac{1}{2} \) | \( \frac{1}{2} \) | \( -\frac{1}{2} \) |
| \( d^e e \) | \( -\frac{1}{2} \) | \( \frac{3}{2} \) | \( -\frac{1}{2} \) |

Table 3. The \( I_{3R} \) and \( B_L \) quantum numbers of the various leptoquark states
plain the HERA anomaly. To see this note that in hadron colliders, we can produce pairs of leptoquarkinos at the same rate as the $t\bar{t}$ pair due to identical color content. Moreover, due to R-parity conservation the leptoquarkino decay leads to a missing energy signal as follows: $\Sigma \rightarrow e^+\bar{u}$ with $\bar{u} \rightarrow u + \chi^0$ or $\Sigma \rightarrow u + e^+$ with $e^+ \rightarrow e^+ + \chi^0$. In both the cases we have $e^+, u$ plus missing energy in the final state. If the $\lambda_{33}$ coupling is comparable to $\lambda_{11}$ as is allowed, then the branching ratio for leptoquarkino decay to electrons will be 50%. This signal is similar to the top signal at the Tevatron with one crucial difference. The dilepton branching ratio from the leptoquark pairs is 100% compared to only 10% from the top pairs. There will be no $\mu^+\mu^-$ events. The branching ratio to $e^+e^-$ will be 25% compared to only 1% from the top. The present observations should therefore lead to lower limits on the mass of the lightest leptoquarkino pairs. Thus an excess of dilepton pairs over that expected from the top productions, or an excess in $e^+e^-$ channel over $\mu^+\mu^-$ will be a clear signal of leptoquark productions at Tevatron. In addition, leptoquarks will also give rise to harder leptons and larger missing energy events. One could also look for the signals of leptoquarkino production in $e^+\gamma$ colliders. The leptoquark or leptoquarkino can be singly produced in such a collider. The scalar leptoquark will be produced in the association of an antiquark. As discussed before in the usual SUSY theories, the leptoquark further decays into a quark and a positron with the final state consisting of positron +jets. The leptoquarkino however will be produced along with a squark. This leptoquarkino will then decay into a lepton and a squark. The final state has electron +jets+missingenergy. The missing energy part will then disentangle the leptoquarkino signal. In the gauge mediated SUSY breaking scenario, the final state has either a hard photon or $\tau^+\tau^-$ along with electron +jets+missingenergy.

Another important implication of our model is that at the renormalizable level, there are any operator which can give a non-negligible charged current signal at HERA without conflicting with $\pi^+ \rightarrow e^+\nu_e$ decay constraint. Thus if a significant amount of charged current like events are observed at HERA this model will have to be further extended.

8. Grandunification with low mass $W_R$

There has been long standing interest in the question of whether low scale supersymmetric left-right models lead to unification of couplings. Several examples of low scale scale left-right models that grand unify already exist in the literature\[44\]. But all of them lead to arbitrary R-parity violation and are therefore outside the philosophical framework of this paper. It is interesting that it is possible to find a low mass $W_R$ SUSYLR model which
leads to grand unification. The basic strategy was to observe that if in a given model, the combination of the one loop beta function coefficients:

\[ \Delta b \equiv 5b_1 - 12b_{2L} + 7b_3 \]  \hspace{1cm} (19)

vanishes, then the new physics scale in this model can be low consistent with the present precision measurements of the gauge couplings. It turns out that if the parity symmetry (but not the $SU(2)_R$) is broken at the GUT scale, then one can choose the following spectrum of particles above the $M_R$ scale to have the $\Delta b = 0$ and thereby a $W_R$ mass in the TeV range: the multiplets are: two bidoublets $\Phi(2, 2, 0)$; the right handed triplets that lead to the see-saw mechanism: $\Delta^c(1, 3 + 2) + \bar{\Delta}^c(1, 3, -2)$, a color octet field neutral under the electroweak gauge group and one $B - L$ neutral $SU(2)_L$ triplet. The one loop evolution equations for the gauge couplings in this case is displayed in Fig.4. The unification scale is safely in the range of $10^{16}$ GeV or so. The important point for our discussion is that this model respects automatic R-parity conservation and also solves the SUSY CP problem as discussed in the text. Whether this model also solves the strong CP problem is not clear yet and is presently under investigation. The unification group is $SO(10)$ or any other group of which $SO(10)$ is a subgroup.

9. Conclusion

In this talk we have shown that the class of supersymmetric left-right models that implement the see-saw mechanism for neutrino masses have two very desirable feature: (i) automatic R-parity conservation and (ii) solution to the susy CP problem due to constraints of parity symmetry. These results are independent of the scale of the left-right symmetry breaking. On the other hand if the $W_R$ scale is in the TeV range, one also has a rather elegant solution to the strong CP problem without the need for an axion. A crucial prediction of this class of models that solve the strong CP problem is that the electric dipole moment of the neutron is of order $10^{-25}$ to $10^{-26}$ emc which is measurable in the next round of proposed experiments. Finally we discuss how the SUSYLR model can incorporate the leptoquarks without the need for R-parity violation in order to explain the HERA high $Q^2$ anomaly.

APPENDIX: Avoiding Sneutrino VEVs

In this Appendix we will show that if in the minimal SUSY LR model one includes non-renormalizable Planck scale induced terms, the ground
state of the theory can be $Q^e$-conserving even for $<\tilde{\nu}^c>=0$. For this purpose, let us briefly recall the argument of Ref. [18]. The part of the potential containing $\tilde{L}^c$, $\Delta^c$ and $\bar{\Delta}^c$ fields only has the form (see Appendix B or [18])

$$V = V_0 + V_D,$$

(20)

where

$$V_0 = \text{Tr}(i\eta^c \tilde{L}^c \tilde{L}^c^T \tau_2 + \mu_{\Delta}^{\tilde{\Delta}})^2 + \mu_3^2 \text{Tr}(\Delta^c \Delta^c^T) + \mu_4^2 \text{Tr}(\bar{\Delta}^c \bar{\Delta}^c^T) + \mu_5^2 \tilde{L}^c \tau_2 \Delta^c \tilde{L}^c,$$

(21)

and

$$V_D = \frac{g^2}{8} \sum_m (\tilde{L}^c^T \tau_m \tilde{L}^c + \text{Tr}(2\Delta^c \tau_m \Delta^c + 2\bar{\Delta}^c \tau_m \bar{\Delta}^c)^2 + \frac{g^2}{8} (\tilde{L}^c \tilde{L}^c - 2 \text{Tr}(\Delta^c \Delta^c - \bar{\Delta}^c \bar{\Delta}^c))^2,$$

(22)

Note that if $<\tilde{\nu}^c>=0$ then the vacuum state for which $\Delta^c = \frac{1}{\sqrt{2}}v\tau_1$ and $\bar{\Delta}^c = \frac{1}{\sqrt{2}}v'\tau_1$ is lower than the vacuum state $\Delta^c = v \left( \begin{array}{cc} 0 & 0 \\ 0 & 1 \end{array} \right)$ and $\bar{\Delta}^c = v' \left( \begin{array}{cc} 0 & 1 \\ 0 & 0 \end{array} \right)$. However, the former is electric charge violating. The only way to have the global minimum conserve electric charge is to have $<\tilde{\nu}^c>\neq 0$. On the other hand, if we have non-renormalizable terms included in the theory, the situation changes: for instance, let us include non-renormalizable gauge invariant terms of the form (inclusion of other non-renormalizable terms simply enlarges the parameter space where our conclusion holds):

$$W_{NR} = \frac{\lambda}{M^2} |\text{Tr}(\Delta^c \tau_m \bar{\Delta}^c)|^2.$$

(23)

This will change $V$ to the form:

$$V = V_0 + V_{NR} + V_D,$$

(24)
where $V_0$ and $V_D$ are given before and $V_{NR}$ is given by

$$V_{NR} = \frac{\lambda \mu}{M} [\text{Tr}(\Delta^c \tau_m \Delta^c)]^2 + \frac{4\lambda \mu \Delta}{M} [\text{Tr}(\Delta^c \tau_m \Delta^c)][\text{Tr}(\Delta^c \tau_m \Delta^c)] + \Delta^c \leftrightarrow \bar{\Delta}^c + \text{etc.} \quad (25)$$

For the charge violating minimum above, this term vanishes but the charge conserving minimum receives a nonzero contribution. Note that the sign of $\lambda$ is arbitrary and therefore, by appropriately choosing $\text{sgn}(\lambda)$ we can make the electric charge conserving vacuum lower than the $Q_{em}$-violating one. In fact, one can argue that, since we expect $v^2 - v'^2 \approx \frac{f^2(M_{SUSY})^2}{16\pi^2}$ in typical Polonyi type models, the charge conserving minimum occurs for

$$f < 4\pi \left(\frac{4\lambda \mu \Delta}{M}ight)^2 \frac{v}{M_{SUSY}}.$$  

For $\lambda \approx 1$, $\mu_\Delta \approx v \approx M_{SUSY} \approx 1\text{TeV}$ and $M = M_P$, we get $f \leq 10^{-3}$ if $v - M_{SUSY}$. We have assumed that the right handed scale is in the TeV range. The constraint on $f$ of course becomes weaker for larger values of $\mu_\Delta$. 

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Figure Caption

**Figure 1:** A typical two loop graph that contributes to the infinite phase of the gluino mass.

**Figure 2:** One loop contribution to $\theta$ from the gluino virtual state. The solid circles represent the squark mass insertions proportional to $h_u^2$.

**Figure 3:** One loop wino contribution to $\theta$ if the wino masses are complex.

**Figure 4:** Gauge coupling unification with the $M_W$ in the TeV range.
Fig. 1
Fig. 3
Fig. 4