\textit{R Symmetry and the $\mu$ Problem}

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Abstract

A natural origin for the $\mu$ and $\mu B$ parameters of weak scale supersymmetric theories is proposed, applicable to any supersymmetry breaking messenger scale between the weak and Planck scales. Although quite general, it requires supersymmetric interactions to respect an $R$ symmetry with definite quantum numbers, and it requires some new scale of symmetry breaking. The required $R$ symmetry distinguishes the Higgs boson from the sneutrino, preserves baryon number in operators of dimension four and five, and contains $R$ parity so that the lightest superpartner is stable. This origin for $\mu$ works for a variety of mediation mechanisms, including gauge mediation, gaugino mediation, and boundary condition breaking of supersymmetry. In any of these mediation schemes, our mechanism leads to a real $B$ parameter, and the supersymmetric $CP$ problem is solved. This $R$ symmetry may naturally arise from supersymmetric theories in higher dimensions.
1. The possibility that nature becomes supersymmetric at the weak scale offers a well motivated and exciting scenario for physics beyond the standard model. It is well motivated because it allows a dynamical generation of the weak scale, and an understanding of why this scale is much less than the Planck scale. Furthermore, it gives rise to a highly successful numerical prediction for gauge coupling unification. It is exciting because weak scale supersymmetry will be thoroughly tested at hadron colliders over the next decade.

Nevertheless, the underlying structure of the fields and interactions of the weak scale supersymmetric theory contains three puzzles:

- In non-supersymmetric field theories there are three distinct types of fields, corresponding to particles with spin 0, 1/2 and 1. In this case, there is no doubt as to what distinguishes the Higgs field from the lepton doublet field. In contrast, in supersymmetric field theories there are just two types of fields: vector multiplets and chiral multiplets. Hence it is now no longer clear what distinguishes a Higgs field, $H$, from a matter field, $M$. In particular, the down type Higgs and lepton doublets have identical gauge and spacetime properties; what distinguishes the Higgs boson from the sneutrino?

- Phenomenologically, interactions cannot be the most general allowed by known gauge and spacetime symmetries. The superpotential must contain interactions of the form $MMH$ for quark and lepton masses, and most probably $MMHH$ at the weak scale for neutrino masses. Yet other forms of superpotential interactions, such as $MH$, $M^3$, and $M^4$ are highly constrained by neutrino masses and proton decay. Some of these interactions are either highly suppressed or forbidden.

- A superpotential interaction of the form $HH$ is a special and intriguing case. On one hand it must be highly suppressed since a coefficient of order the Planck scale or unified mass scale would remove the Higgs doublets from the low energy theory. On the other hand it cannot vanish since otherwise there is a massless charged fermion coming from the Higgs/vector multiplets. Indeed the theory is only realistic if the coefficient, $\mu$, is of order the weak scale, leading to the well known $\mu$ problem. Why should this supersymmetric parameter be of order the supersymmetry breaking scale?

The other great mystery of low energy supersymmetry is the origin of supersymmetry breaking. Like the supersymmetric interactions, a great deal about the structure of the supersymmetry breaking interactions is governed by the requirement of consistency with experiment. However, nothing determines the “messenger” scale, $M_m$: the highest scale at which the supersymmetry breaking interactions of squarks, sleptons, Higgs and gauginos are local. In supergravity theories this locality is maintained up to the Planck scale [1], but other methods of transporting supersymmetry breaking to the standard model superpartners, such as gauge mediation [2, 3], gaugino mediation [4], and boundary condition supersymmetry breaking [5],
have messenger scales that can be anywhere between the weak and Planck scales.

In this paper we study the consequences of imposing an $R$ symmetry on the supersymmetric interactions of effective theories with weak scale supersymmetry. This symmetry will allow us to address all three of the puzzles listed above in the context of supersymmetry breaking with an arbitrary messenger scale. We will describe a continuous symmetry, $U(1)_R$, although a discrete subgroup is sufficient for our purposes. The soft supersymmetry breaking operators break the $R$ symmetry, since they include Majorana gaugino masses, and therefore holomorphic, $A$ term, scalar interactions. Using this $R$ symmetry, we find a new mechanism for solving the $\mu$ problem for arbitrary messenger scales, and this requires a unique choice for the $R$ quantum numbers of matter and Higgs fields. Furthermore, we find that this is also the unique choice which accounts for the absence of $MH$, $M^3$ and $M^4$ superpotential interactions, while allowing consistency with quark-lepton unification. Finally, this $R$ symmetry forces a distinction between Higgs and lepton doublets.

There are two well-known classes of solutions to the $\mu$ problem. One class corresponds to modifying physics at the Planck scale by adding non-renormalizable operators [6]; the other changes the physics at the weak scale, as is the case in the Next-to-Minimal Supersymmetric Standard Model (NMSSM) [7]. The first approach requires supersymmetry breaking to be mediated at the Planck scale, while the second requires a departure in the weak scale theory from the Minimal Supersymmetric Standard Model (MSSM). Neither of these requirements apply to our mechanism.

In our mechanism, we suppose that all supersymmetric interactions respect some global symmetry, $G$, which forbids $\mu$ and commutes with both gauge and flavor symmetries. To allow for the possibility of unification of the matter, we require the $G$ charges of all matter multiplets to be equal. We consider that some field acquires a vacuum expectation value (vev) at scale $\Lambda$ between the weak and Planck scales. Supersymmetry breaking interactions, mediated at any scale above $\Lambda$, break $G$ and cause a deformation in the vacuum resulting in the generation of $\mu$ of order the supersymmetry breaking scale. Such theories have been constructed for the case of mediation at the Planck scale [8]. However, for such high mediation scales $\mu$ is not a problem since the Giudice-Masiero mechanism is available. We study the case of an arbitrary mediation scale. Related solutions for the $\mu$ problem in the context of gauge mediation have been proposed [9], but differ from our mechanism in the origin of the vacuum deformation.

2. If we require that our solution to the $\mu$ problem apply to arbitrary mediation scales, we find that the operator giving rise to $\mu$ is uniquely determined. We assume this operator must be present at tree level — if $\mu$ arises radiatively then the soft $\mu BH_u H_d$ term is typically generated at too high a level. Furthermore, $\mu$ must arise from a renormalizable operator. If the operator had a coefficient suppressed by powers of the Planck scale, it would not be possible to get a $\mu$ parameter of the desired size for arbitrary values of the messenger scale. Thus the
operator generating $\mu$ is unique:

$$W_\mu = \lambda X H_u H_d,$$

where $X$ is a standard model gauge singlet chiral superfield. Our mechanism requires that the supersymmetric interactions break some symmetry at a scale $\Lambda$, giving a mass to $X$ and forcing $\langle A_X \rangle = \langle F_X \rangle = 0$. Here $A_X$ and $F_X$ represent the lowest and highest components of the chiral superfield $X$, respectively. Non-zero values for $\langle A_X \rangle$ and $\langle F_X \rangle$ are generated by supersymmetry breaking.

We discover that our global symmetry, $G$, must be an $R$ symmetry from the following argument. The order of magnitude of $\langle A_X \rangle$ and $\langle F_X \rangle$ generated after supersymmetry breaking can be understood from the $G$ charges of both $X$ and the soft supersymmetry breaking operators. If $G$ is a non-$R$ symmetry then $A_X$ and $F_X$ have the same $G$ transformation, so that both $\mu$ and $\mu B$ are generated at the same order in supersymmetry breaking. Hence $B$ is of order $\Lambda$ and much too large: our mechanism requires $G$ to be an $R$ symmetry.

Now, under the $R$ symmetry, there are only two types of supersymmetry breaking terms appearing in the scalar potential. There are holomorphic terms, which we denote by $A$ and assign $R$ charge $-2$, and there are non-holomorphic terms, denoted $m^2$, which have $R$ charge zero. A successful solution to the $\mu$ problem requires $\mu$, and therefore $\langle A_X \rangle$, to be linear in supersymmetry breaking. This requires that $X$ has $R$ charge $+2$ or $-2$, so that $\langle A_X \rangle$ can be generated proportional to $A^*$ or $A$. Interestingly, this automatically leads to $\langle F_X \rangle$ of exactly the right order, since now $F_X$ has $R$ charge 0 or $-4$ and is naturally generated at second order in supersymmetry breaking proportional to $|A|^2$, $m^2$ or $A^2$.

Let us consider our two possible $R$-charge assignments separately. If $X$ has $R$ charge $-2$, then $H_u H_d$ has total $R$ charge $+4$. Due to the presumed existence of unification, all matter fields carry the same $R$ charge. The Yukawa couplings then force equal $R$ quantum numbers for the two Higgs doublets, $R(H_u) = R(H_d)$, so that $R(H) = 2$ and $R(M) = 0$. However, in this case the superpotential interaction $MH$ is allowed, which will push some MSSM matter fields to have masses of order the Planck scale. Therefore, we prefer to consider the charge assignment $R(X) = 2$. Again, the Yukawa couplings require $R(H_u) = R(H_d)$, so we have

$$R(H) = 0 \quad R(M) = 1.$$  \hfill (2)

This $R$ symmetry is extremely powerful: as well as distinguishing between the lepton and Higgs doublet and forbidding $MH$, it also forbids $M^3$ and $M^4$, leading to baryon and lepton number conservation from operators of dimension four and five. Therefore, our solution to the $\mu$ problem has forced us to forbid dangerous dimension four and five operators that might lead to too rapid proton decay. Moreover, if the soft supersymmetry breaking operators provide the only source of $R$ breaking, then $R$ parity remains unbroken, leading to stability of the lightest
With a mild assumption about the origin of neutrino masses, a completely independent argument will lead us to an identical conclusion for the global symmetry $G$. The Yukawa interactions, $M^2 H$, possess a non-$R$ Peccei-Quinn symmetry (PQ: $M(1), H(-2)$) as well as the $R$ symmetry of Eq. (2). In fact, $G$ must be a linear combination of $R$ and PQ, or one of its subgroups. The existence of small neutrino masses strongly suggests that the superpotential also contains $M^2 H^2$. Provided that this interaction is not generated by supersymmetry breaking, this immediately implies that $H$ is neutral under $G$, and hence $G$ must be the $R$ symmetry given in Eq. (2).

3. The simplest model which realizes the above general mechanism for generating the $\mu$ term is given by the superpotential

$$W_0 = f X(Y^2 - \Lambda^2),$$

where $R(X) = 2$ and $R(Y) = 0$. Here we imagine that the scale $\Lambda$ is much larger than the weak scale, although this is not necessary for our mechanism to work. We also impose a discrete symmetry $Y \to -Y$, so that the gauge hierarchy is not destabilized by the generation of a large tadpole operator for a singlet field. Without supersymmetry breaking, the minimum of the potential lies at $\langle A_X \rangle = 0$ and $\langle A_Y \rangle = \Lambda$, satisfying $\langle F_X \rangle = \langle F_Y \rangle = 0$. There is no flat direction at this level, and all the fields have masses of order $\Lambda$.

When we add supersymmetry breaking terms, with a scale $\tilde{m}$ of order the weak scale, the vevs will shift. The most general soft supersymmetry breaking terms are given by

$$\mathcal{L}_{\text{soft},0} = -m_X^2 |X|^2 - m_Y^2 |Y|^2 - (a_f XY^2 - a_\Lambda \Lambda^2 X + \text{h.c.}).$$

Here, $R(m_X^2) = R(m_Y^2) = 0$ and $R(a_f) = R(a_\Lambda) = -2$, and we have used $X$ and $Y$ to denote the scalar fields of the respective chiral superfields. By minimizing the scalar potential, we obtain $\langle X \rangle \simeq (a_\Lambda^* - a_f^*)/4|f|^2 \sim \tilde{m}$ and $\langle F_X \rangle \simeq [(a_\Lambda + a_f)(a_\Lambda^* - a_f^*)/4|f|^2 + m_Y^2]/2f \sim \tilde{m}^2$. The vevs of $X$ and $F_X$ are both of the order of the weak scale as indicated by the previous general analysis. Therefore, if we introduce couplings to the Higgs doublet $W = W_0 + \lambda X H_u H_d$ and $\mathcal{L}_{\text{soft}} = \mathcal{L}_{\text{soft},0} - m_{H_u}^2 |H_u|^2 - m_{H_d}^2 |H_d|^2 - (a_\lambda X H_u H_d + \text{h.c.})$, $\mu$ and $\mu B$ terms of order $\tilde{m}$ and $\tilde{m}^2$ are generated as

$$\mu = \lambda \langle X \rangle \simeq \frac{\lambda (a_\lambda^* - a_f^*)}{4|f|^2},$$

$^1$ $R$ parity forbids the generation of dimension four baryon and lepton number violating operators, even after supersymmetry is broken. While dimension five proton decay operators could be generated with small coefficients proportional to supersymmetry breaking, they are phenomenologically irrelevant.

$^2$ We assume that violations of global symmetries by nonperturbative gravitational effects are sufficiently suppressed.
\[ \mu B = -\lambda \langle F_X \rangle + a_\lambda \langle X \rangle \simeq \frac{(2fa_\lambda - \lambda a_f)(a_\lambda^* - a_f^*)}{8f|f|^2} - \frac{\lambda m_\lambda^2}{2f}, \] (6)

where \( \mu \) and \( \mu B \) are defined by \( W = \mu H_u H_d \) and \( \mathcal{L} = -\mu B H_u H_d + \text{h.c.} \). Although \( \mu = 0 \) for \( a_\lambda = a_f \), it is natural to expect \( a_\lambda \neq a_f \) since the two parameters run differently under renormalization group evolution; a simple realistic example will be given later. An important point here is that there exists a parameter region where the additional Higgs coupling in Eq. (1) does not change the gross dynamics at the high scale. For instance, if \( m_{H_u}^2 \) and \( m_{H_d}^2 \) are sufficiently large, the vev of order \( \Lambda \) is still entirely contained in the \( Y \) field, and the vevs for the Higgs doublets stay smaller than the weak scale. Note that here the various supersymmetry breaking parameters are evaluated at the scale \( \Lambda \). Thus, both \( m_{H_u}^2 \) and \( m_{H_d}^2 \) can be positive without conflicting with electroweak symmetry breaking. Below the scale \( \Lambda \), the heavy fields \( X \) and \( Y \) are integrated out and the effective theory contains only the Higgs doublets with the \( \mu \) and \( \mu B \) parameters given by Eqs. (5, 6). The soft supersymmetry breaking parameters must be further evolved from \( \Lambda \) to \( \tilde{\Lambda} \) using the renormalization group equations of this effective theory (MSSM), to evaluate electroweak symmetry breaking.

An important requirement for our mechanism is that \( \Lambda \) must be smaller than the messenger scale of supersymmetry breaking, \( \Lambda < M_m \). Therefore, the superpotential Eq. (3) itself is not sufficient for a complete solution of the \( \mu \) problem (except for the supergravity mediation case), since we have introduced by hand a mass parameter, \( \Lambda \), smaller than the fundamental scale. A complete solution, however, is obtained if we generate the scale \( \Lambda \) by the dynamics of strong gauge interactions. Consider, for example, the \( SU(2)_S \) gauge theory with four doublet chiral superfields \( Q_i \) \( (i = 1, \ldots, 4) \) with the following superpotential:

\[ W_{0,\text{tree}} = fX \left( Y^2 - (QQ) \right) + f'X^a(QQ)_a. \] (7)

This superpotential explicitly breaks a flavor \( SU(4)_F \) symmetry of the \( Q_i \) down to \( SP(4)_F; (QQ) \) and \( (QQ)_a \) \( (a = 1, \ldots, 5) \) denote singlet and five-dimensional representations of \( SP(4)_F \) given by suitable combinations of gauge invariants \( Q_i Q_j \). The strong dynamics of the \( SU(2)_S \) gauge theory is described by the effective superpotential

\[ W_{0,\text{eff}} = W_{0,\text{tree}} + S \left( (QQ)^2 + (QQ)_a^2 - \Lambda^4 \right), \] (8)

where \( S \) is an additional Lagrange multiplier chiral superfield \([12]\). For a relatively large value of the coupling \( f' \), the vacuum lies at \( (QQ) = \Lambda^2 \) and \( (QQ)_a = 0 \), so that the superpotential \( W_{0,\text{eff}} \) is effectively reduced to Eq. (3). Note that the original tree-level superpotential, Eq. (7), does not contain any mass parameters and is invariant under the \( U(1)_R \) symmetry with \( R(X) = R(X^a) = 2 \) and \( R(Q_i) = 0 \). In fact, it is the most general superpotential consistent with the combined \( R \) and \( SP(4)_F \) symmetries. (A linear term in \( X \) is forbidden either by requiring that the superpotential not contain any mass parameters, or by imposing an anomalous discrete \( Z_3 \)
symmetry under which all the fields are transformed by \( \exp(2\pi i/3) \). It is also important that \( U(1)_R \) does not have an anomaly for \( SU(2)_S \) (i.e. \( \Lambda \) does not carry \( U(1)_R \) charge), so that the previous general argument is not affected by the strong \( SU(2) \) gauge dynamics.

We now consider an application of our mechanism to realistic theories. We find that the mechanism fits beautifully into the framework where small neutrino masses are generated by integrating out right-handed neutrino fields through the see-saw mechanism \[13\]. We consider the following theory. In addition to the usual three generations of standard-model quark and lepton superfields, \( Q, U, D, L \) and \( E \), we introduce three right-handed neutrino superfields \( N \).

Here, we have omitted generation indices. The Yukawa couplings are given by

\[
W_{\text{Yukawa}} = y_u Q U H_u + y_d Q D H_d + y_e L E H_d + y_\nu L N H_u. \quad (9)
\]

We also introduce the \( U(1)_X \) gauge symmetry, contained in \( SO(10)/SU(5) \), under which various fields transform as \( Q(1), U(1), D(-3), L(-3), E(1) \) and \( N(5) \). This gauge symmetry is broken by the vevs of the fields \( \Phi(10) \) and \( \bar{\Phi}(-10) \) through the superpotential

\[
W_{\text{Breaking}} = f X (\Phi \bar{\Phi} - (QQ)) + f' X^a (QQ)_a. \quad (10)
\]

Here \( (QQ) \) and \( (QQ)_a \) are gauge invariants consisting of \( Q_i \), the doublets under the strong \( SU(2)_S \) gauge interaction (see discussion around Eqs. (7, 8)). Note that the above superpotentials, Eqs. (9, 10), do not contain any mass parameters and are invariant under the \( U(1)_R \) symmetry,

\[
R(Q) = R(U) = R(D) = R(L) = R(E) = R(N) = 1, \quad R(H_u) = R(H_d) = R(\Phi) = R(\bar{\Phi}) = R(Q_i) = 0 \quad \text{and} \quad R(X) = R(X^a) = 2.
\]

For a relatively large \( f' \), the dynamics of the \( SU(2)_S \) gauge interaction cause the condensation of \( (QQ) = \Lambda^2 \), which is transmitted to the vevs of the fields \( \Phi(10) \) and \( \bar{\Phi}(-10) \) through the superpotential

\[
W_{\text{Breaking}} = f X (\Phi \bar{\Phi} - (QQ)) + f' X^a (QQ)_a. \quad (10)
\]

The vevs for all the other fields are zero at this stage:

\[
\langle X \rangle = \langle X^a \rangle = \langle (QQ)_a \rangle = 0.
\]

After introducing soft supersymmetry breaking operators, the vevs of the fields shift. In particular, non-vanishing vevs for \( X \) and \( F_X \) are generated as \( \langle X \rangle \sim \tilde{m} \) and \( \langle F_X \rangle \sim \tilde{m}^2 \), as long as holomorphic soft supersymmetry breaking parameters are not subject to the special relation \( a_{X \Phi \bar{\Phi}} = a_{X(QQ)} \) at the scale \( \Lambda \). In fact, it is quite natural to expect that the \( A \) terms for \( X \Phi \bar{\Phi} \) and \( X(QQ) \) are different since they are renormalized differently above the scale \( \Lambda \); for example, they receive contributions from \( U(1)_X \) and \( SU(2)_S \) gauginos, respectively. Therefore, by introducing the couplings

\[
W_{\text{Masses}} = \kappa \Phi N^2 + \lambda X H_u H_d, \quad (11)
\]

the Majorana masses \( M_R \) for the right-handed neutrinos of order \( M_R \sim \langle \Phi \rangle \sim \Lambda \) and \( \mu \) and \( \mu B \) parameters of order \( \mu \sim B \sim \tilde{m} \) are generated. As in the previous example, the superpotential

\[
W = W_{\text{Yukawa}} + W_{\text{Breaking}} + W_{\text{Masses}}
\]

is the most general renormalizable superpotential consistent
with the gauge $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_X \times SU(2)_S$ and global $U(1)_R \times SP(4)_F$ symmetries of the theory, after removing a linear term in $X$ as before.

We finally discuss how our general mechanism works explicitly in various supersymmetry breaking scenarios. In gauge mediated supersymmetry breaking, our mechanism requires that the mass of the messenger fields, $M_m$, is larger than $\Lambda$. Since the holomorphic supersymmetry breaking terms ($A$ terms) required for the $\mu$-term generation are small at the messenger scale $M_m$, they must be generated by renormalization group evolution from $M_m$ to $\Lambda$. This can be accomplished, for example, by giving non-trivial $SU(2)_S$ or $U(1)_X$ quantum numbers to the messenger fields. In the case of gaugino mediation and boundary condition supersymmetry breaking, our mechanism requires that the compactification scale is larger than $\Lambda$. In these cases, the relevant $A$ terms of order the weak scale may already exist at the compactification scale, so we do not necessarily have to rely on renormalization group evolution for their generation. In any of these mediation mechanisms, $A$ is real in the basis where the gaugino masses are real (except for the case of gaugino mediation with tree-level $A$ terms), and our origin for $\mu$ and $\mu B$ then leads to a real $B$ parameter: the supersymmetric $CP$ problem is solved.

4. In this paper we have proposed an origin for the parameters $\mu$ and $\mu B$ of the minimal supersymmetric standard model, which is applicable for any messenger scale, $M_m$. Although quite general, it does require specific symmetries and interactions. Both $\mu$ and $\mu B$ parameters arise from the superpotential interaction $XH_uH_d$. A stage of symmetry breaking occurs at some scale $\Lambda < M_m$, giving a mass of order $\Lambda$ to $X$, while determining $\langle A_X \rangle = \langle F_X \rangle = 0$. Providing the form for the superpotential is guaranteed by an $R$ symmetry, with the quantum numbers of Eq. (2), the soft supersymmetry breaking operators, with coefficients $A$ and $m^2$, lead to a small readjustment of the vacuum, giving

$$\mu \approx A^*, \quad \mu B \approx |A|^2, m^2.$$  (12)

This $R$ symmetry provides a distinction between Higgs and matter superfields, and forbids superpotential interactions that would otherwise lead to baryon number violation at too rapid a rate. Although $R$ is broken by supersymmetry breaking, the discrete $R$ parity survives so that the lightest superpartner is stable. In the case that the original $R$ symmetry is continuous, an $R$ axion will be produced by the underlying dynamics which breaks supersymmetry. If all the $R$ breaking effects are generated spontaneously (including the constant term in the superpotential needed to cancel the cosmological constant), the dominant mass contribution to the $R$ axion will come from the QCD anomaly of the $R$ symmetry. In this case, the $R$ axion provides a solution to the strong $CP$ problem [15].

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3 In anomaly mediation [14], the holomorphic supersymmetry breaking parameter associated with the scale $\Lambda$ is large — of the order of the gravitino mass $\sim 10$ TeV. Thus $\lambda \sim 10^{-2}$ is needed to generate a $\mu$ parameter of the correct size. Then, to avoid a too large $\mu B$ term, a cancellation between two contributions, such as $f a$ and $\lambda a\Lambda$ in Eq. (3), is required at the $1\%$ level.
At first sight our $R$ symmetry appears to be in conflict with grand unification: since $R$ forbids $H_uH_d$, it also forbids the corresponding mass term for the colored Higgs triplets of unified theories. However, this turns out to be a virtue — such mass terms need to be forbidden to avoid too large a proton decay rate mediated by triplet Higgsino exchange. The colored partners of $H_u,d$ must become heavy by acquiring mass terms coupling them to other colored states of the theory. This occurs in the missing partner and Dimopoulos-Wilczek mechanisms; however, although these mechanisms are consistent with an underlying $U(1)_R$ symmetry, in the simplest such models $U(1)_R$ is broken at the unification scale, so our $\mu$ generation mechanism may not work in these cases. In contrast, in Kaluza-Klein grand unification the desired colored Higgs mass terms arise while preserving $U(1)_R$ symmetry, so that our $\mu$ generation mechanism works well in this case.

The $U(1)_R$ symmetry is so crucial in providing an understanding of the form for the interactions in the superpotential, it is important to seek its origin. Higher dimensional theories are particularly interesting since they have an enlarged set of supersymmetry transformations, which results in a global $R$ symmetry in the equivalent four dimensional description. In the case of a five dimensional grand unified theory, compactification breaks the unified gauge symmetry and also the $SU(2)_R$ symmetry to $U(1)_R$, so that precisely the $R$ charges considered here may arise.

**Acknowledgements**

Y.N. thanks the Miller Institute for Basic Research in Science for financial support. This work was supported in part by the Director, Office of Science, Office of High Energy and Nuclear Physics, of the U.S. Department of Energy under Contract DE-AC03-76SF00098, and in part by the National Science Foundation under grant PHY-00-98840.
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