Relaxation of the Dynamical Gluino Phase and Unambiguous Electric Dipole Moments

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Abstract

We propose a new axionic solution of the strong CP problem with a Peccei-Quinn mechanism using the gluino rather than quarks. The spontaneous breaking of this new global U(1) at $10^{11}$ GeV also generates the supersymmetry breaking scale of 1 TeV (solving the so-called $\mu$ problem at the same time) and results in the MSSM (Minimal Supersymmetric Standard Model) with R parity conservation. In this framework, electric dipole moments become calculable without ambiguity.
CP nonconservation is a fundamental issue in particle physics. We know that CP is not conserved in $K - \bar{K}$ mixing (i.e. $\epsilon \neq 0$)[1] and in $K$ decay (i.e. $\epsilon' \neq 0$)[2]. However, only an upper limit $(0.63 \times 10^{-25} \text{e}\cdot\text{cm})$ exists for the electric dipole moment (edm) of the neutron[3]. This may not be so bothersome until we realize that the currently accepted theory of strong interactions, i.e. quantum chromodynamics (QCD), actually violates CP through the instanton-induced term[4]

$$\mathcal{L}_\theta = \theta_{QCD} \frac{g_s^2}{32\pi^2} \epsilon_{\mu\nu\lambda\beta} G^\mu_a G^\nu_a$$

(1)

where $g_s$ is the strong coupling constant, and

$$G_a^{\mu\nu} = \partial^{\mu} G_a^{\nu} - \partial^{\nu} G_a^{\mu} + g_s f_{abc} G_b^{\mu} G_c^{\nu}$$

(2)

is the gluonic field strength. The value of $\theta_{QCD}$ must then be less than $10^{-10}$ in magnitude (instead of the expected order of unity) to account for[5] the nonobservation of the neutron edm. This is known as the strong CP problem.

Another fundamental issue in particle physics is supersymmetry. It allows us to solve the hierarchy problem so that our effective theory at the electroweak energy scale ($M_W$) is protected against large radiative corrections. However, this requires the scale of soft supersymmetry breaking ($M_{SUSY}$) to be not much higher than $M_W$. There is no theoretical understanding of why the two scales must be related in this way.

In the following we address the question of how the strong CP problem is to be solved in the context of supersymmetry. We find that it can be achieved with a new kind of axion[6, 7] which couples to the gluino rather than to quarks. In a natural implementation of this idea, we find that the breaking of supersymmetry must have the same origin as the axion. The scale of electroweak symmetry breaking is also related. This works because the so-called $\mu$ problem is being solved along the way.

In our framework, the $\theta_{QCD}$ contribution to any quark edm is canceled exactly by the minimization of the dynamical gluino phase. Hence the calculation of edm’s in the MSSM (Minimal Supersymmetric Standard Model) becomes unambiguous. The possibility of cancellation[8] among other different CP nonconserving contributions[9] to edm’s can now be pursued without fear of contradiction.
With the addition of colored fermions, the parameter \( \theta_{QCD} \) of Eq. (1) is replaced by
\[
\bar{\theta} = \theta_{QCD} - \text{Arg Det } M_u M_d - 3 \text{ Arg } M_{\tilde{g}},
\] (3)

where \( M_u \) and \( M_d \) are the respective mass matrices of the charge = 2/3 and charge = -1/3 quarks, and \( M_{\tilde{g}} \) is the mass of the gluino. The famous Peccei-Quinn solution\[6\] is to introduce a dynamical phase to the quark masses which then relaxes \( \bar{\theta} \) to zero. The specific realization of this requires an axion\[7\] which is ruled out experimentally\[10\]. Two other axionic solutions have been proposed which are at present consistent with all observations. The DFSZ solution\[11\] introduces a singlet scalar field as the source of the axion but its mixing with the doublet scalar fields which couple to the quarks is very much suppressed. The KSVZ solution\[12\] introduces new heavy quarks so that the axion does not even couple directly to the usual quarks. Neither scheme requires supersymmetry.

In the context of supersymmetry however, it is clear that the simplest and most natural thing to do is to attach the axion to the gluino rather than to the quarks in Eq. (3). Because of the structure of supersymmetry, all other superparticles will be similarly affected. This is then a very strong hint that it may have something to do with the breaking of supersymmetry. As shown below with our proposed singlet complex scalar field \( S \), whose phase contains the axion, all soft supersymmetry breaking parameters are of order \( |\langle S \rangle|^2/M_{Pl} \), where \( M_{Pl} \sim 10^{19} \text{ GeV} \) is the Planck mass. Hence a value of \( 10^{11} \text{ GeV} \) for \( |\langle S \rangle| \), which is allowed by astrophysical and cosmological constraints, would imply \( M_{\text{SUSY}} \sim 1 \text{ TeV} \). Since the electroweak symmetry breaking terms are also among this group, it does not require any stretch of the imagination to find \( M_W \) and \( M_{\text{SUSY}} \) to be only an order of magnitude apart.

It is known\[13\] that a continuous global \( U(1)_R \) symmetry\[14\] can be defined for the MSSM. The quark (\( \tilde{Q}, \tilde{u}^c, \tilde{d}^c \)) and lepton (\( \tilde{L}, \tilde{e}^c \)) chiral superfields have \( R = +1 \) whereas the Higgs (\( \tilde{H}_u, \tilde{H}_d \)) chiral superfields and the vector superfields have \( R = 0 \). The superpotential
\[
\hat{W} = \mu \hat{H}_u \cdot \hat{H}_d + h_u \hat{Q} \cdot \hat{H}_u \tilde{u}^c + h_d \hat{Q} \cdot \hat{H}_d \tilde{d}^c + h_e \hat{L} \cdot \hat{H}_d \tilde{e}^c
\] (4)
has \( R = +2 \) except for the \( \mu \) term (which has \( R = 0 \)). In the above, the Yukawa couplings \( h_{u,d,e} \) are nonhermitian matrices in flavor space. The
resulting Lagrangian is then invariant only with respect to the usual discrete
$R$ parity, i.e.
\[ R \equiv (-1)^{3B+L+2J}, \]  
where $B$ is baryon number, $L$ lepton number, and $J$ spin angular momentum,
hence $R$ is even for particles and odd for superparticles.

We now propose to make $U(1)_R$ an exact global symmetry of the supersymmetric
Lagrangian, as well as that of all the supersymmetric breaking
terms. We introduce the composite operator
\[ \mu(\hat{S}) \equiv \frac{1}{M_{Pl}} (\hat{S})^2, \]  
where the singlet superfield $\hat{S}$ has $R = +1$. Our model is then defined by
the new superpotential
\[ \hat{W} = \mu(\hat{S}) \hat{H}_u \cdot \hat{H}_d + m_s^2 \mu(\hat{S}) 
+ h_u \hat{Q} \cdot \hat{H}_u \hat{u}^c + h_d \hat{Q} \cdot \hat{H}_d \hat{d}^c + h_e \hat{\tilde{L}} \cdot \hat{H}_d \hat{e}^c, \]  
which has $R = +2$, thus yielding a supersymmetric Lagrangian which is
invariant under $U(1)_R$, together with the following set of supersymmetry
breaking terms which are also invariant under $U(1)_R$:
\[ \Delta \mathcal{L} = |\mu(\hat{S})|^2 \left[ \hat{Q}^\dagger Y_Q \hat{Q} + \hat{u}^c \hat{Y}_u \hat{u}^c + \hat{d}^c \hat{Y}_d \hat{d}^c + \hat{\tilde{L}}^\dagger \hat{Y}_L \hat{\tilde{L}} + \hat{\tilde{e}}^c \hat{Y}_{\tilde{e}} \hat{\tilde{e}}^c \right] 
+ \left\{ \mu(\hat{S})^\dagger [k_u \hat{Q} \cdot \hat{H}_u \hat{u}^c + k_d \hat{Q} \cdot \hat{H}_d \hat{d}^c + k_e \hat{\tilde{L}} \cdot \hat{H}_d \hat{e}^c] + h.c. \right\} 
+ |\mu(\hat{S})|^2 [y_u |H_u|^2 + y_d |H_d|^2 + (k_u H_u \cdot H_d + h.c.)] 
+ \left\{ \mu(\hat{S})^\dagger [k_3 \tilde{\lambda}_3^a \tilde{\lambda}_3^a + k_2 \tilde{\lambda}_2^a \tilde{\lambda}_2^a + k_1 \tilde{\lambda}_1 \tilde{\lambda}_1] + h.c. \right\}, \]  
where $\tilde{\lambda}_3^a$ is the gluino octet, $\tilde{\lambda}_2^a$ the $SU(2)_L$ gaugino triplet, and $\tilde{\lambda}_1$ the
$U(1)_Y$ gaugino singlet. The parameters $k_{1,2,3}$ and $\mu$ are complex, whereas
$y_{u,d}$ are real. The matrices $Y_{Q,L}$ and $Y_{u,d,e}$ are hermitian, whereas $k_{u,d,e}$ are
nonhermitian. Obviously, we have assumed in the above that the source of
all supersymmetry breaking terms is $\mu(\hat{S})$. Together with $U(1)_R$, this solves
the so-called $\mu$ problem in the MSSM, because the scale of $\mu(\hat{S})$ is $|\langle S \rangle|^2/M_{Pl}$
which is of order 1 TeV for $|\langle S \rangle| \sim 10^{11}$ GeV, instead of the typical grand
unification scale of $10^{16}$ GeV.

The Lagrangian $\Delta \mathcal{L}$ describes the interaction of $\mu(\hat{S})$ with sfermions,
Higgs doublets and gauginos only. However, a complete description of our
model requires the self interactions of the singlet to be specified as well. The pure singlet contribution \( m^2_s \mu(S) \) in \( W \) is allowed by the symmetries of the model and the mass parameter \( m^2_s \) is \textit{a priori} arbitrary. Through the \( F \)-term contributions, this induces a positive mass-squared parameter for the singlet: \( m^2_F \equiv 4m^4_s/M^2_{Pl} \). However, interactions at higher energies at or near the Planck scale can provide an additional mass-squared parameter \( m^2_0 \) as well as a quartic coupling \( m^2_s \). Through the \( F \)-term contributions, this induces a positive mass-squared parameter for the singlet: \( m^2_F + m^2_0 \). Since the Higgs doublets have vanishing \( R \) charges, the electroweak breaking cannot have any effect on the fate of \( U(1)_R \). The only way to break it is to allow the singlet to develop a nonvanishing vacuum expectation value. This can happen only when \( M^2_s < 0 \) so that \( v^2_s = -M^2_s/2\lambda_s \). Since \( m^2_F \) is positive, \( m^2_0 \) should be negative enough to induce a negative \( M^2_s \). This implies that \( m^2_s \) cannot be as large as \( M^2_{Pl} \) as it would leave \( U(1)_R \) unbroken; hence \( |m^2_s| \sim |m^2_0| \sim v^2_s \) is a natural choice. The singlet field could then be expanded around \( v_s \) as

\[
S(x) = [v_s + s(x)] e^{i\phi(x)},
\]

where \( \phi(x) \) is the corresponding Nambu–Goldstone boson[15] which has a strictly flat potential, and \( s(x) \) is a real scalar field with a mass of order \( v_s \).

It is clear from the above that our \( U(1)_R \) plays the role of what is usually called \( U(1)_{PQ} \)[6]. Whereas the conventional \( U(1)_{PQ} \) applies to the usual quarks and leptons, our \( U(1)_{PQ} \) applies only to the superparticles, and the gluino is the only colored fermion having a nonvanishing \( U(1)_{PQ} \) current:

\[
J^5_{\mu} = \overline{\lambda}^a \gamma_\mu \gamma^5 \lambda^a,
\]

where we have used the four-component notation: \( \lambda^a = (\tilde{\lambda}_3^a, \tilde{\lambda}_3^a) \). Now the gluino also contributes to the color current with respect to which \( J^5_{\mu} \) has a nonvanishing quantum anomaly:

\[
\partial^\mu J^5_{\mu} = \frac{6g^2_s}{32\pi^2} \epsilon^{\mu\nu\alpha\beta} G^\mu_a G^\nu_a G^\alpha\beta \quad .
\]

Since \( J^5_{\mu} \) couples to \( \phi \) as \( \partial^\mu \phi J^5_{\mu} \), the effective QCD vacuum angle takes the form

\[
\tilde{\theta} = \theta_{QCD} + 6\phi(x),
\]

where the nondynamical phases in the quark mass matrices and the phase of the complex constant \( k_3 \) can be included in \( \theta_{QCD} \) by a chiral rotation. In close
analogy with the KSVZ scenario[12], our $\varphi(x)$ also receives a potential from
the instanton background so as to develop a vacuum expectation value which
enforces $\vartheta \equiv 0$ [i.e. $\langle \varphi \rangle = -\theta_{QCD}/6$], to all orders in perturbation theory.
Rather than the quarks, it is thus the gluino which realizes the Peccei-Quinn
mechanism of solving the strong CP problem.

The axion, $a \equiv v_s[\varphi(x) - \langle \varphi \rangle]$, has a mass and lifetime given by

$$m_a \sim m_\pi \frac{f_\pi}{f_a} , \quad \tau(a \to 2\gamma) \sim \left( \frac{m_\pi}{m_a} \right)^5 \tau(\pi \to 2\gamma) ,$$

where its decay constant $f_a$ is equal to $v_s/6$. Our axion is not a DFSZ
axion[11] as it does not couple to quarks and leptons; it is also not a KSVZ
axion[12] as it does not couple to unknown colored fermion multiplets beyond
the MSSM spectrum. We may call it the gluino axion[16] as it is induced
by promoting the masses of the gauginos to local operators.

Let us choose $v_s \sim 10^{11}$ GeV, which is in the middle of the range allowed
by astrophysical and cosmological bounds[17] on $f_a$. The effective theory
below $v_s$ is then a replica of the MSSM with the effective $\mu$ parameter

$$\mu_{eff} = \frac{v_s^2}{M_{Pl}} e^{-i\theta_{QCD}/3} \sim 10^3 \text{ GeV} \times e^{-i\theta_{QCD}/3},$$

which is the right scale for supersymmetry breaking. This seesaw mechanism
for the $\mu$ parameter results from the introduction of the composite operator
given by Eq. (6) into the theory, the dynamics of which are presumably
dictated by physics at or near the Planck scale. On the other hand, the scale
of the $\mu$ parameter is fixed by the astrophysical and cosmological bounds on
the axion decay constant $f_a$ which gives (or receives) a meaning to (from)
the intermediate scale $v_s$.

The low-energy effective theory is the softly broken MSSM with $R$ parity
conservation. Indeed, after replacing the effective $\mu$ parameter [Eq. (14)] and
$\langle \varphi \rangle = -\theta_{QCD}/6$ into the effective Lagrangian [Eq. (8)], we obtain

$$\mathcal{L}_{MSSM}^{soft} = \hat{Q}^\dagger M_{Q}^{2} \hat{Q} + \bar{u}^\dagger M_{u}^{2} \bar{u} + \bar{d}^\dagger M_{d}^{2} \bar{d} + \bar{c}^\dagger M_{c}^{2} \bar{c} + \bar{L}^\dagger M_{L}^{2} \bar{L} + e^\dagger M_{e}^{2} e +$$

$$+ \left\{ A_u \hat{Q} \cdot H_u \bar{u} + A_d \hat{Q} \cdot H_d \bar{d} + A_c \bar{L} \cdot H_d \bar{c} \right\} + h.c. \} \right) + M_{H_d}^{2} |H_d|^2 + (\mu_{eff} BH_u \cdot H_d + h.c.) + \left\{ M_3 \lambda_3^\dagger \lambda_3 + M_2 \lambda_2^\dagger \lambda_2 + M_1 \lambda_1^\dagger \lambda_1 + h.c. \right\} ,$$

(15)

which is nothing but the soft supersymmetry-breaking part of the MSSM
Lagrangian. It is in fact this part of the Lagrangian that possesses all sources
of CP violation through the complex $A$ parameters, the gaugino masses, and $\mu_{\text{eff}}$ itself. The explicit expressions for the mass parameters in Eq. (15) read as follows. The gaugino masses are given by

$$M_3 = |k_3|\mu_{\text{eff}}^* , \quad M_2 = k_2\mu_{\text{eff}}^* , \quad M_1 = k_1\mu_{\text{eff}}^* ,$$

which are not necessarily universal. The soft masses for the Higgs sector are given by

$$M_{H_u}^2 = y_u|\mu_{\text{eff}}^2| , \quad M_{H_d}^2 = y_d|\mu_{\text{eff}}^2| , \quad \mu_{\text{eff}}B = |\mu_{\text{eff}}|^2\left(\frac{4m_s^2}{v_s^2} + k_\mu\right) ,$$

and are responsible for electroweak symmetry breaking, with similar expressions for the mass-squared matrices of the sfermion fields. As discussed before, $m_s^2/v_s^2 \sim \mathcal{O}(1)$ so that $B$ parameter is also at the weak scale. Finally the $A$ parameters are given by

$$A_u = \mu_{\text{eff}}^*k_u , \quad A_d = \mu_{\text{eff}}^*k_d , \quad A_e = \mu_{\text{eff}}^*k_e ,$$

which do not have to be proportional to $h_u$, $h_d$, and $h_e$ of Eq. (7) as in the constrained MSSM.

As noted before, all mass scales of the MSSM Lagrangian [Eq. (15)] are fixed in terms of $|\mu_{\text{eff}}|$. More than this, the phase of $\mu_{\text{eff}}$, i.e. $-\theta_{\text{QCD}}/3$, contributes universally to all mass parameters which are complex. However, the phases of the gaugino masses as well as those of the $A$ and $B$ terms also depend on the $k$ parameters. Hence if the flavor structure of these matrices is not the same as those of the usual quarks and leptons, then the CP violation in flavor-changing processes is a powerful probe[18] into this sector of the effective theory. In the calculation of electric dipole moments due to supersymmetry[9], these CP phases can be considered as they are without worrying about whether there is an additional contribution from $\bar{\theta}$.

In conclusion, we have presented in the above a simultaneous solution to two hierarchy problems, i.e. why $\bar{\theta}$ is so small (the strong CP problem) and why $\mu$ is 1 TeV and not $10^{16}$ GeV (the $\mu$ problem), as well as the related issue of why $M_{W}$ and $M_{SUSY}$ are only one order of magnitude apart. The primary difference between our approach and previous other attempts[19] lies in the fact that the gaugino masses are promoted here to local operators given by $\mu(\hat{S})$. Indeed, finite bare mass terms for the gauginos would automatically break the $U(1)_{PQ}$ symmetry, making it impossible for the relaxation of $\bar{\theta}$ to
zero. As it is, \( \langle S \rangle \) serves two important purposes: its magnitude determines the scale of supersymmetry breaking and its phase solves the strong CP problem.

Let us summarize our proposal.

(i) We work in the framework of supersymmetry and identify \( U(1)_{PQ} \) as \( U(1)_R \) which contains the usual \( R \) parity as a discrete subgroup.

(ii) We require the supersymmetric Lagrangian and all supersymmetry breaking terms to be invariant under \( U(1)_R \).

(iii) We implement this with the composite operator \( \mu(S) \equiv (\hat{S})^2/M_{Pl} \) where \( \hat{S} \) is a singlet superfield having \( R = +1 \).

(iv) The spontaneous breaking of \( U(1)_R \) generates an axion and relaxes the effective QCD vacuum angle \( \vartheta \) to zero, using the dynamical gluino phase, thus solving the strong CP problem.

(v) The existing astrophysical and cosmological bounds on the axion decay constant implies a supersymmetry breaking scale of 1 TeV.

(vi) The effective Lagrangian at low energy is that of the MSSM with \( R \) parity conservation. All mass scales are of order 1 TeV, thus solving the \( \mu \) problem and the related issue of why \( M_W \) and \( M_{SUSY} \) are only an order of magnitude apart.

(vii) Since \( \vartheta = 0 \) in this consistent supersymmetric theory, electric dipole moments can be calculated unambiguously from the other explicit CP violating terms of the MSSM.

D.A.D. acknowledges the hospitality of the UCR Physics Department, where this work was initiated. The research of E.M. was supported in part by the U. S. Department of Energy under Grant No. DE-FG03-94ER40837.
References


[14] Under $U(1)_R$, the scalar components of a chiral superfield transforms as $\phi \rightarrow e^{i\theta R}\phi$, whereas the fermionic components transform as $\psi \rightarrow e^{i\theta(R-1)}/\psi$. 


[16] Details such as the precise $a \to \gamma \gamma$ coupling and other phenomenological implications will be discussed in a forthcoming paper: D. A. Demir and E. Ma, in preparation.

