Hidden sector gaugino condensation and the model-independent axion

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(Article begins on next page)
Hidden Sector Gaugino Condensation and the Model-independent Axion

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In the effective field theory framework, we consider the effect of supersymmetry breaking via gaugino condensation and supergravity in the hidden sector gauge group on the hidden sector vacuum angle $\theta_h$. The $\theta_h$ parameter dependence of the potential yields phenomenologically acceptable invisible axion solutions if the $U(1)_R$ symmetry is broken down to a discrete subgroup $\mathbb{Z}_N$ with $N \geq 3$ ($N = 4$ is marginal). Anomalous $U(1)$ superstring models are good candidates for this invisible axion resolution of the strong CP puzzle.
The QCD vacuum angle $\theta$ is bounded very strongly by the upper bound on the neutron electric dipole moment to $\theta < 10^{-9}$. The almost vanishing strong interaction parameter has led to the so-called strong CP puzzle. The axion may be the most attractive solution of the strong CP puzzle, though other possibilities such as the massless up quark or calculable may be viable.

However, the Peccei-Quinn (PQ) global symmetry, the global symmetry needed for the axion solution, or the up quark chiral symmetry is ad hoc in gauge theories, since the gravitational interaction is expected to break all global symmetries explicitly. Even if the axion solution is attractive theoretically and cosmologically, one has to overcome the puzzle of global symmetry in the presence of the gravitational interaction. Because superstring theory may provide a consistent approach to quantum gravity, it seems reasonable to require that any global symmetry that is introduced be obtained in string theory. In this regard, the discovery of the model-independent axion (MIA) in superstring models is of most fundamental importance. This MIA suffers from the axion decay constant problem.

On the other hand, this axion decay constant problem can be circumvented with anomalous $U(1)$ gauge groups. The argument of is worth repeating. Some 4D superstring models have an anomalous $U(1)$, but the anomaly is not troublesome due to the Green-Schwarz term which effectively introduces a coupling of the MIA with the anomalous $U(1)$ gauge boson

$$
\epsilon_{MNOPQRSTUV} B^{MN} \text{Tr} F^{OP} \langle F^{QR} \rangle \text{Tr} \langle F^{ST} \rangle \langle F^{UV} \rangle
$$

where the vacuum expectation value is of order the compactification scale. In 4D, the coupling gives a quadratic term $\sim M_c \partial^2 a_{MI} A_\mu$, showing that the MIA becomes the longitudinal component of the anomalous $U(1)$ gauge boson $A_\mu$. This breaks the anomalous $U(1)$ gauge symmetry and the PQ symmetry associated with translation of $a_{MI}$, leaving a global symmetry below the compactification scale.

In more detail, the way this works is as follows. The original nonlinearly realized tree level global symmetry is the symmetry,

$$
a_{MI} \rightarrow a_{MI} + \alpha F_{MI} \\
\psi_i \rightarrow e^{g_i q_i A_{MI} / M_c} \psi_i
$$

where $\alpha$ is a constant, $F_{MI}$ is the MIA decay constant and $\psi_i$ are the matter fields that transform nontrivially under the symmetry. Because of ii, the MIA has only derivative interactions.

The anomalous $U(1)$ gauge symmetry takes the form

$$
A_\mu \rightarrow A_\mu + (1/g_1) \partial_\mu \beta(x) \\
\psi_i \rightarrow e^{i g_i \beta(x)} \psi_i \\
a_{MI} \rightarrow a_{MI} + M_c \beta(x)
$$

where $M_c$ is the compactification scale, so that covariant derivatives have the form

$$
D_\mu \psi_i = \left( \partial_\mu - ig_i q_i A_\mu \right) \psi_i \\
D_\mu a_{MI} = \partial_\mu a_{MI} - g_1 M_c A_\mu
$$

Below the scale $M_c$, the MIA and its interactions completely disappear, because it is gauged away to become the longitudinal component of the anomalous $U(1)$ gauge boson. The heavy $A_\mu$ fields disappear from the low energy theory, and the original global PQ symmetry is irrelevant at low energies. However, there is an anomalous global symmetry in the low energy theory, which is a linear combination of the PQ symmetry and a global anomalous gauge symmetry, with $\beta(x) = -\alpha F_{MI} / M_c$.

Thus, below the compactification scale $M_c$, the MIA does not exist, but there is a global $U(1)$ symmetry which we will call $U(1)_X$. The $U(1)_X$ is an anomalous global symmetry because it involves the anomalous $U(1)$. This makes it an appropriate candidate to be the PQ symmetry in the low energy theory, below $M_c$.

In spontaneously broken gauge models, it has long been known that the symmetry breaking of a $U(1)$ gauge and a $U(1)$ global symmetry by the VEV of one Higgs scalar can preserve a global symmetry below the spontaneous symmetry breaking scale. The discussion above is in essence this ’t Hooft mechanism.

In models in which SUSY is broken by gaugino condensation in a hidden sector and supergravity, we expect this condensation to spontaneously break the $U(1)_X$ at a scale $F_X \sim 10^{12}$ GeV. When the $U(1)_X$ breaks spontaneously, a new pseudoGoldstone boson is produced which we will call $a_X$. It is an axion, because the $U(1)_X$ symmetry is...
anomalous. We emphasize that it is not the MIA — it is “made of” the hidden sector gaugino fields that condense to break $U(1)_X$ at the scale $F_X$.

If QCD were the only confining force, $a_X$ would be the perfect choice for the invisible axion in superstring models. However, we are assuming that the hidden sector that is required for supersymmetry breaking involves another confining interaction. Thus we have two vacuum angles — $\theta$ of QCD and $\theta_h$ of the hidden sector confining gauge group. The $a_X$ contributes to both angles:

$$\theta = \frac{a_X}{F_X} + \theta^0, \quad \theta_h = \frac{a_X}{F_X} + \theta_h^0$$

where $\theta^0$ and $\theta_h^0$ are parameters given by the CP violations of weak interactions and hidden sector physics. In general, $\theta_h^0 \neq \theta^0$. Since the hidden sector scale is much larger than the QCD scale, we might naively expect the hidden sector interactions to give a much larger contribution than QCD to the potential. These hidden sector interactions are minimized for $\theta_h = 0$, and if they dominate, we will find $a_X \simeq -F_X \theta_h^0$. But then the QCD angle $\theta$ is set to $\theta^0 - \theta_h^0$, not to zero, and the strong CP puzzle is not resolved. To set both $\theta$ and $\theta_h$ to zero, we would need two independent axions. But it seems that except for the global symmetry related to the MIA $U(1)_{MI}$, there is no bosonic global symmetry in superstring models.

Thus to implement the invisible axion with a confining hidden sector, we must find theories in which the $\theta_h$ dependence of the vacuum energy is suppressed. In fact, some suppression is built into supersymmetric theories of this kind because the bare gaugino mass is much smaller than the scale $F_X$ of gaugino condensation, because it arises only due to the supersymmetry breaking produced by the supergravitational interactions. However, this is not enough to suppress the $\theta_h$ dependence to acceptable levels.

If the hidden sector gluino carried an unbroken anomalous $U(1)$ symmetry like the PQ symmetry in a model with a massless $u$ quark, that would eliminate the $\theta_h$ dependence of the vacuum energy entirely and allow the QCD interactions to determine the vacuum value of $a_X$ and set $\theta = 0$. A $U(1)_R$ symmetry rotates the hidden sector gaugino fields and is therefore a candidate for such a symmetry. However, we do not expect any such global continuous symmetry to be exact in the presence of gravitational interactions. Thus the best we can do is to find models in which the contribution of the hidden sector to the $a_X$ potential is sufficiently suppressed, because of the precise form of the $U(1)_R$ breaking interactions. For this purpose, we look for models with the following properties:

(i) A $Z_N$ subgroup of the continuous $U(1)_R$ symmetry survives even in the presence of the gravitational effects that explicitly break the continuous $U(1)_R$

(ii) The spontaneous $U(1)_R$ violation at the scale $F_X$ occurs only through hidden sector gaugino condensation and supergravity.

Let us now examine the consequence of these assumptions. Viable superstring models would lead to effective low energy supersymmetry with nonrenormalizable terms suppressed by string or Planck scale. The dominant low energy physics from a D=4 supergravity theory is characterized by three functions, the superpotential, the Kähler potential, and the gauge kinetic function,

$$W(\phi_i), \quad K(\phi_i, \phi_i^*), \quad f_i^a(\phi_i)$$

where $\{\phi_i\}$ is a set of chiral superfields, and the subscript $a$ of the gauge kinetic function runs over different gauge groups, $G = \prod_a G_a$. Supersymmetry is broken when $\langle G^i \rangle \neq 0$ where $G = K + \log |W|^2$ and $G^i = \partial G/\partial \phi_i$ (from here on, we use gravitational units in which the Planck mass is one). The gravitino mass becomes $m_{3/2} = e^{G/2} = e^{K/2}|W|$.

In fact, both the superpotential and

$$f_i^a(\phi_i) W^{a\alpha} W_\alpha$$

---

1The translation of this effective field theory statement into the language of the full field theory below the string scale is this. Because there are contributions to the breaking of the two symmetries from both of the two scales, $M_c$ and $F_X$, there are two Goldstone fields, $a_{MI}$ and the “invisible axion”, $a$, which is a linear combination of phases of $U(1)_X$ breaking singlet fields. One combination of $a_{MI}$ and $a$, $a_X \simeq a_{MI} + e a$ where $e \sim F_a/M_c$, is eaten to become the longitudinal degree of the anomalous gauge boson. The low energy invisible axion is the orthogonal linear combination, $a_X \simeq a - e a_{MI}$. It contains a small component of the model-independent axion.

2And, of course, it would ensure that the mass term for the hidden sector gaugino remained zero even after SUSY breaking.

3This may be possible in string theory even if there are no continuous global symmetries.
(where $W^a_\alpha$ is the gauge field strength superfield for the group $a$) depend only on left-handed superfields, and their $F$ terms appear in the Lagrangian. But if gaugino condensation occurs at a very large scale, we must also consider similar terms in the Lagrangian with higher powers of $W^a_\alpha W^a_\alpha$, where $a$ refers to the hidden sector gauge group, because below the condensation scale, the powers of $W^a_\alpha W^a_\alpha$ are replaced by powers of $\Lambda_h^4$, where $\Lambda_h$ is the large hidden sector scale. These terms can be parameterized as follows:

$$\int d^2\theta \sum_{a,n} f_n^a(\phi_i)(W^a_\alpha W^a_\alpha)^n$$

(8)

where the $f_n^a$ functions are holomorphic in $\phi$. Because of our assumption (ii), these are the terms that give the dominant contribution to $R$-symmetry breaking in the low energy theory.

Furthermore, because the spontaneous $R$ symmetry breaking in the hidden sector is assumed to arise from gaugino condensation only, the $R$ breaking terms from (8) in the low energy theory are evaluated at zero VEV of chiral fields (assumption (ii) again). We can take

$$f_n^a = \epsilon_n^a S$$

(9)

where $S$ is a chiral superfield. Terms with more powers of chiral fields, such as $S^l(l > 1), S A^n B^m \cdots (n, m \geq 1)$, etc. do not contribute since we assumed VEV’s of chiral fields are vanishing. By power counting, the $\epsilon$ parameters in (1) are suppressed by appropriate powers of the Planck mass, $M_P$. In general, there could be a whole set of $S$ fields, but the effect of one will be enough to illustrate our mechanism.

From (8) and (9), we get contributions to the potential energy of the form

$$V = \left[ (\lambda \cdot \lambda) + \sum_{n \geq 2} \epsilon_n^a (\lambda \cdot \lambda)^n \right]^2$$

(10)

where the $\lambda$s are the hidden sector gaugino fields and where we have normalized the coefficient of the first term

$$\epsilon_1^a = 1.$$ 

(11)

If only the first term exists, i.e. $f_n^a = 0$ for $n \geq 2$, then (10) is not $R$ violating.¶

Let us now consider the effect of our assumption that a $Z_N$ subgroup of $U(1)_R$ remains unbroken. Under such a $Z_N$ subgroup,

$$\lambda \rightarrow e^{2\pi i/N} \lambda$$

(12)

Thus if $N$ is even, $(\lambda \lambda)^{2N/2}$ is invariant, for any integer $j$, and if $N$ is odd $(\lambda \lambda)^{jN}$ is invariant. In either case, the sum over $n$ in (10) will be restricted. The most important $R$-violating term in the potential will result from the first allowed term in the sum over $n$, and will have the form

$$\epsilon_m(\lambda \cdot \lambda)(\lambda \cdot \lambda)^m + \text{h.c.}$$

(13)

where $m$ is the lowest integer greater than 1 for which $\epsilon_m \neq 0$, which is $N + 1$ for $N$ odd, and $N + 1$ for $N$ even. The interesting values of $m$ are $m \geq 3$. The $m = 2$ ($N = 2$) case is not interesting because in that case, the $\Sigma_2$ symmetry does not forbid the gaugino mass term.

Notice that the $Z_N$ symmetry must be a subgroup of the continuous $R$ symmetry because the gauge kinetic term, $\int d^2\theta W^a_\alpha W^a_\alpha$ must be invariant — the gauginos and the gauge bosons in the multiplet must transform differently.

Our question is what is the effect of the above form of $R$ symmetry breaking on the dependence on the hidden sector parameter $\theta_h$. To consider this problem, we take an effective field theory viewpoint below the hidden sector scale $\Lambda_h$. The light degree of freedom arising from the gaugino condensation is the pseudoscalar meson $\eta'$ whose mass is about $\Lambda_h$. The parameters describing the effective theory are $\epsilon_m, \theta_h$, and the scale for the gaugino condensation.

---

4Because are primarily interested in gaugino condensation, we consider only $f$’s that are gauge singlets — otherwise there would be more general terms here, as below in (8).

5Another way to see this is that we could choose the $R$ quantum number of $S$ to preserve $R$ invariance.
Since this phenomenon is very similar to the chiral symmetry breaking of QCD, we first recall some basic facts about the $\theta$ dependence in QCD. As the simplest example, let us consider up quark condensation in one-flavor QCD. The mass term in this theory is given by

$$\mathcal{L}_{mass} = -m_u \bar{u}u$$

Formally, we can assign the following $U(1)$ chiral transformation,

$$u \rightarrow e^{i\alpha} u, \quad \bar{u} \rightarrow e^{i\alpha} \bar{u}, \quad m \rightarrow e^{-2i\alpha} m, \quad \theta \rightarrow \theta + 2\alpha$$

From the above chiral symmetry, we expect the following effective potential below the chiral symmetry breaking scale,

$$V = \frac{1}{2} m_u \Lambda^3 e^{i\theta} - \frac{1}{2} \lambda_1 v^3 e^{i\frac{\theta}{v}} - \frac{1}{2} \lambda_2 m_u v^3 e^{i\frac{\theta}{v}} + \lambda_3 m_u^2 \Lambda^2 e^{2i\theta} + \lambda_4 \frac{\rho^6}{\Lambda^6} e^{2i\frac{\vartheta}{v}} - 2i\theta + \cdots + \text{h.c.}$$

where $\cdots$ denotes higher order terms, $\lambda$'s are couplings of order 1, $\langle \bar{u}u \rangle = v^3 e^{i\eta/v}$, and the QCD scale $\Lambda$ is inserted to make up the correct dimension. In addition, $e^{\pm i\theta}, e^{\pm 2i\theta}$ etc is multiplied to respect the $U(1)$ symmetry. Note that if $m_u \neq 0$ and $\theta$ is not a dynamical variable, then the strong CP puzzle is not solved. However, if $m_u = 0$ then only the $m_u$-independent terms survive, leading to

$$V = -\frac{1}{2} \lambda_1 v^3 e^{i\frac{\theta}{v}} + \lambda_4 \frac{\rho^6}{\Lambda^6} e^{2i\frac{\vartheta}{v}} - 2i\theta + \cdots + \text{h.c.}$$

Thus, redefining the $\eta$ field as $\eta'$

$$\eta' = \eta - v\theta,$$

the $\theta$ dependence is completely removed from $V$. The $\theta$ parameter is unphysical if a quark is massless. As is well known the massless up quark scenario solves the strong CP puzzle even though it obtains a constituent quark mass when the chiral symmetry is broken. Of course, as we discussed in the introduction, the relevance of this solution to the strong CP puzzle hinges on its viability in hadron physics phenomenology [3].

To compare with the gaugino condensation case, we comment on the possible violation of the chiral symmetry by gravitational interactions for the case of the massless up quark. Even if the mass term is not present, chiral symmetry breaking nonrenormalizable interactions may be generated,

$$\frac{1}{M_P} \bar{u}u\sigma\bar{\sigma}e^{-2i\theta}, \frac{1}{M_P} \bar{u}u\sigma\sigma e^{-i\theta}, \text{ etc.}$$

where $\sigma$ is the singlet field carrying no chiral charge. The first term will shift $\theta$ by a tiny amount $|\langle \bar{u}u \rangle|^{2/3}/M_P^2 \sim 10^{-38}$ which is ridiculously small. But if some scalar singlets develop a large expectation value, one expects a sizable $\theta \sim |\langle \sigma \rangle|^2/M_P^2 |\langle \bar{u}u \rangle|^{1/3}$. The upper bound on $\theta$ implies $|\langle \sigma \rangle| < 10^4$ GeV.

In the same way, the spontaneous $R$ symmetry breaking by the hidden sector gaugino condensation dictates the $\theta_h$ parameter dependence of the potential through the explicit $R$ breaking $\epsilon_m$ dependence. As in QCD, we can study the form of the potential by treating $\epsilon_m$, as a spurion field, and imposing the symmetry

$$\lambda \rightarrow e^{i\alpha} \lambda, \quad \epsilon_m \rightarrow e^{-2m\alpha} \epsilon_m, \quad \theta_h \rightarrow \theta_h + 2\ell(G)\alpha$$

$^6$Note that we are not trying to keep careful track of the powers of $2\pi$ here, so the distinction between the scale $\Lambda$ and the amplitude for $\eta$ production $v$ has not been made very precisely.
where $\ell(G)$ is the index of the adjoint representation of the gauge group $G$. We assume the hidden sector group as $SU(K)$ for which $\ell = K$. The effective interaction is

$$V_{\text{eff}} = \frac{1}{2} \bar{\lambda} \lambda \lambda + \epsilon_m \bar{\lambda} \lambda (\lambda \lambda)^m + \frac{\Lambda_h^4}{\Lambda_h^{4K}} (\lambda \lambda)^K e^{-i\theta_h} + \cdots + \text{h.c.}$$

(21)

which includes the terms

$$\left( \frac{v}{M_p} \right)^{3m-1} v^4 e^{i(m-1)\eta/v} + \left( \frac{v}{\Lambda_h} \right)^{3K} \Lambda_h^4 e^{i(\Delta - \theta_h)} + \cdots + \text{h.c.}$$

(22)

where we have replaced $\epsilon_m$ by the appropriate number of powers of $M_p$. Redefining $\eta' = \eta - v\theta_h/K$ (because the second term in (22) is much larger than the first and enforces this identification), we obtain the $\theta_h$ dependence

$$\frac{v^{3m+3}}{M_p^{3m-1}} e^{-i(m-1)(\theta_h/K + \eta'/v)} + \cdots + \text{h.c.}$$

(23)

This is the dominant term depending on $\theta_h$.

Let us now discuss the $\theta_h$ dependence in (23), remembering that $m$ is the lowest integer greater than 1 for which $\epsilon_m \neq 0$. Clearly, the larger $m$, the smaller the $\theta_h$ dependence. Even if the hidden sector scale is large, $\Lambda_h \sim 10^{12}$ GeV, the height of the $\theta_h$ potential is very low. The situation is summarized in table 1.

For $N = 3$ the suppression is adequate over almost all of the range. $N = 4$ is marginal. But all higher $N$ provide adequate suppression.

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<td>N=3</td>
<td>m=4</td>
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<td>N=4</td>
<td>m=3</td>
<td>$\sim 10^{-8} - 10^{10}$</td>
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<tr>
<td>N=5</td>
<td>m=6</td>
<td>$\sim 10^{-50} - 10^{-26}$</td>
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Table 1: The $\theta_h$ dependence of potential in GeV$^4$ units for $\Lambda_h = 10^{12} - 10^{13}$ GeV and Planck mass suppression factor of $10^{19} - 2.4 \times 10^{18}$ GeV.

In conclusion, we have studied the implication of the hidden sector gaugino condensation on breaking of $Z_N$ subgroups of $U(1)_R$. Under the assumption that $U(1)_R$ breaking VEV’s of chiral fields do not exist, we showed that anomalous $U(1)$ models with $N \geq 3$ ($N = 4$ is marginal) are viable candidates containing phenomenologically acceptable invisible axion models.

In string models, discrete symmetries like the $Z_N$ we posit arise naturally. In fact, the discrete symmetries were used before [14] to obtain an accidental PQ symmetry. Our idea differs from Ref. [14] in that in our case the PQ symmetry is present in anomalous $U(1)$ gauge models and we employ a specific discrete symmetry, i.e. the discrete subgroup of $U(1)_R$, to suppress the hidden sector contributions to the axion mass. In Ref. [14], the PQ symmetry arises accidentally and the extra confining group is not considered.

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