Strong CP-Problem in Superstring Theory

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Abstract
We apply the solution for the strong CP-problem in the 4-dimensional superstring theory recently proposed by Ibáñez and Lüst to Calabi-Yau type models and study its phenomenological aspects. In Calabi-Yau type models there seem to be phenomenologically difficult problems in the axion decoupling from the neutral gauge currents and the compatibility between the proton stability and the cosmological bound on the axion. DFSZ type invisible axion mechanism which works without heavy extra colored fields may be more promising than KSVZ axion in the viewpoint of proton stability.
Strong CP problem\cite{1} still now remains being the obscure situation in superstring theory. It is not known whether PQ mechanism\cite{2} works or not in it. The pseudoscalar fields which come from the antisymmetric tensor are found not to work as the invisible axion against the initial expectation\cite{3}. The introduction of the global $U(1)$ symmetry to the superstring is also difficult\cite{4}. Recently Ibáñez and Lüst pointed out that the target space modular invariant 4-dimensional superstring theory, especially, $(0, 2)$ orbifolds has the very interesting properties for the strong CP problem\cite{5}. Their findings are following: 1) PQ symmetry is automatically built in the theory as the Kähler transformation associated to the target space duality. 2) The soft supersymmetry breaking terms are real and then there is no extra dangerous contribution to the EDM of neutrion in contrast with the ordinary supersymmetric theory. In this note we apply their solution to Calabi-Yau type models and examine their phenomenological aspects in detail.

The target space modular invariance\cite{6} is very useful to know the low energy effective Lagrangian of superstring\cite{7} \cite{8}. This invariance is known to be kept in all order of string perturbation and is usually expected to be retained even in the non-perturbative effect. If we define the one of Kähler moduli as $T = R^2 + iD$ ( $R$ is the overall radius of the compactified manifold and $D$ is the model dependent axion originating from the antisymmetric tensor), the target space modular transformation is defined as

$$T \rightarrow T' = \frac{aT - ib}{icT + d}$$

(1)

where $ad - bc = 1$ and $a, b, c, d$ are integers.\footnote{If there are more Kähler moduli fields as the general Calabi-Yau case, the transformation becomes more complicated one. But (1) is expected to be contained in it. In the following study we confine ourselves to the one Kähler modulus case, for simplicity.} Under this transformation, the every chiral fields $\phi_i$ have modular weight $n_i$ and transform as

$$\phi_i \rightarrow \phi'_i = (icT + d)^{n_i} \phi_i.$$ 

(2)

The generalized Kähler potential $G = K(\Phi_i, \Phi_i^*) + \log |W(\Phi_i)|^2$ is known to be invariant under the following Kähler transformation

$$K(\Phi_i, \Phi_i^*) \rightarrow K(\Phi_i, \Phi_i^*) + F(\Phi_i) + F^*(\Phi_i^*),$$

$$W(\Phi_i) \rightarrow \exp(-F(\Phi_i))W(\Phi_i),$$

(3)

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(3)
where $\Phi_i$ represents $T$ and matter fields $\phi_i$. The target space modular transformation (1) induces a Kähler transformation with $F(T) = \log(i c T + d)^3$. The superpotential $W(T, \phi_i)$ of the charged matter fields $\phi_i$ is generally written as

$$W(T, \phi_i) = \sum_{i,j,k} \lambda_{ijk}(T) \phi_i \phi_j \phi_k.$$  \hspace{1cm} (4)

From eqs. (2) and (3), it is required that $\lambda_{ijk}(T)$ should transform in the following way,

$$\lambda_{ijk}(T) \to (icT + d)^{-3-n_i-n_j-n_k} \lambda_{ijk}(T).$$  \hspace{1cm} (5)

Here we should note that in order for the transformation (3) to be the symmetry of the theory, the fermion fields $\psi_I$ with the canonically normalized kinetic term should be also transformed as

$$\psi_I \to \exp[-\frac{1}{4} q_I (F - F^*)] \psi_I,$$  \hspace{1cm} (6)

where $q_I = 1$ for gauginos and $q_I = -1 - \frac{2}{3} n_I$ for matter fermions ($n_I$ is the modular weight)\[5\] [9]. As a result, this symmetry has generally a triangle anomaly

$$\frac{i}{32\pi^2} \sum_I q_I \text{Tr} T_I^2 (F - F^*) F^a_{\mu\nu} \tilde{F}^{a\mu\nu}.$$  \hspace{1cm} (7)

$T_I$ is the generator of the gauge group in the representation of the fermion $\psi_I$. This means that if there is a scalar field which has a flat potential except for the effect of the color anomaly, an invisible axion mechanism can work due to this symmetry. In this mechanism the PQ-like symmetry is built in the effective Lagrangian of the superstring as its own property from the beginning.

Now we apply these mechanisms to Calabi-Yau type superstring models. In these models the low energy effective Lagrangian is composed of the various charged matter fields. These models have the following $E_6$ charged fields

$$Q(3, 2)_1/3, 1/\sqrt{\sigma}, -1/\sqrt{\sigma}; \quad \bar{u}(3^*, 1)_{-4/3, 1/\sqrt{\sigma}, -1/\sqrt{\sigma};} \quad \bar{d}(3^*, 1)_{2/3, 1/\sqrt{\sigma}, 3/\sqrt{\sigma};}$$

$$\ell(1, 2)_{-1, 1/\sqrt{\sigma}, 3/\sqrt{\sigma};} \quad S(1, 1)_{0, 1/\sqrt{\sigma}, -5/\sqrt{\sigma};} \quad \bar{e}(1, 1)_{-2, 1/\sqrt{\sigma}, -1/\sqrt{\sigma};}$$

$$h(1, 2)_{-1, -2/\sqrt{\sigma}, 2/\sqrt{\sigma};} \quad h(1, 2)_{-1, -2/\sqrt{\sigma}, -2/\sqrt{\sigma};} \quad S(2, 1)_{0, 4/\sqrt{\sigma}, 0;}$$

$$g(3, 1)_{-2, 3, -2/\sqrt{\sigma}, 2/\sqrt{\sigma};} \quad \bar{g}(3^*, 1)_{2, 3, -2/\sqrt{\sigma}, -2/\sqrt{\sigma};}$$

and also $E_6$ singlet fields. The above representations of each fields stand for the ones under $SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)_\psi \times U(1)_\chi (\subset E_6)$. Without loss of generality the

4We adopt $E_6$ model as the starting point of our following arguments. The generalization to the models with more extra $U(1)$’s is straightforward.
superpotential $W$ for these fields is expressed

$$W = \lambda_1 Qg + \lambda_2 Q\bar{u}h + \lambda_3 Q \bar{d}h' + \lambda_4 Q\bar{g}\ell + \lambda_5 \bar{u}\bar{d} + \lambda_6 \bar{u}\bar{g}\bar{e}$$

$$+ \lambda_7 \bar{d}g S_1 + \lambda_8 \bar{g}g S_2 + \lambda_9 hh' S_2 + \lambda_{10} \ell h S_1 + \lambda_{11} \ell h' \bar{e}.$$  \hspace{1cm} (9)

We abbreviated the generation indices. $\lambda_2, \lambda_3, \lambda_7$ and $\lambda_8$ terms are relevant to $\bar{\theta} = \theta + \arg \det M$ through the mass matrix $M$ of the colored fields. As seen from eqs.(5) and (7), under the target space modular transformation (1) parameter $\theta$ shifts through these couplings

$$\theta \rightarrow \theta + \log \det \left( \frac{\lambda_8(T)S_2}{\lambda_8^*(T)S_2^*} \right) + \log \det \left( \frac{\lambda_2(T)h}{\lambda_2^*(T)h^*} \right) + \log \det \left( \frac{\lambda_3(T)h'}{\lambda_3^*(T)h'^*} \right)$$

$$= \theta + \sum_{j,k} (-3 - n_j - n_k) \log \left( \frac{icT + d}{-icT^* + d} \right).$$  \hspace{1cm} (10)

The summation should be taken over the colored fields contributing to these couplings.

To realize the symmetry breaking of the standard model in our considering models, all of $S_1, S_2, h, h'$ must have VEVs. This suggests that the linear combination of the phase of $S_1, S_2, h$ and $h'$ will work as an axion.\footnote{In general $\lambda_7$ term is phenomenologically unfavorable and we assume it to be zero in eq.(10) and the following arguments. Particularly, $S_2$ should have a VEV at an intermediate mass scale $\sim 10^{11}$GeV.}

In the most superstring theories it is known that there exist the extra color triplets $g, \bar{g}$ which couple to the singlet fields similar to $S_2$ as $\lambda_8$ term. Ibáñez and Lüst pointed out that the phase part of the Kähler transformation associated to the target space modular transformation (1) plays the $U(1)_{PQ}$ role through $\bar{g}g S_2$ coupling like the invisible axion model of KSVZ\cite{KSVZ}. As seen from the above arguments we should note that the axion in this mechanism has the following properties:

(i) ordinary quarks, leptons and doublet Higgs scalars have this $U(1)_{PQ}$ charge as DFSZ axion model\cite{DFSZ};

(ii) singlet field $S_2$ decouples from ordinary quarks and leptons not due to $U(1)_{PQ}$ but the extra gauge symmetry,

(iii) extra triplets $g$ and $\bar{g}$ generally couple not only to singlet $S_2$ but also to ordinary quarks and leptons through $\lambda_1, \lambda_4, \lambda_5$ and $\lambda_6$ terms.
These properties yield the non-trivial phenomenological problems on this mechanism at least in the Calabi-Yau type models. That is, (i) makes the axion the mixture of doublet Higgs $h, h'$ and singlets $S_1, S_2$. As a result there appears non-trivial axion decoupling problem from the neutral gauge currents in Calabi-Yau type models. Here we should note that the extra gauge symmetry also plays an important role to guarantee the flatness of the scalar potential of $S_2$ (i.e. the absence of $S_2^3$ term in $W$) other than (ii). (iii) brings the compatibility problem of the proton stability and the cosmological bound on the axion. In the following, we study these problems in detail.

We start from the study of the axion decoupling from the neutral gauge currents in this scenario. At first, we briefly review the generalized axion model [1]. Let’s consider the models which have a set of scalar fields $\phi_i$ with non-trivial PQ charge $\Gamma_i$. We define

$$\phi_i = \frac{1}{\sqrt{2}}(v_i + \eta_i) \exp(i \frac{\xi_i}{v_i}) \quad (11)$$

where $v_i$ is the vacuum expectation value (VEV). The axion field is written as

$$a = \frac{1}{f_a} \sum_i \Gamma_i v_i \xi_i. \quad (12)$$

$f_a$ is the axion decay constant and expressed by

$$f_a^2 = \sum_i \Gamma_i^2 v_i^2. \quad (13)$$

We should note that $f_a$ is dependent on the absolute value of the PQ charge. Only if $\Gamma_i = O(1)$, $f_a$ is the measure of the scale of symmetry breaking. When the symmetry breakings occur, would-be Nambu-Goldstone bosons are absorbed by the gauge bosons. The axion must be orthogonal to these neutral gauge currents $j_\mu^a$,

$$\langle 0 | j_\mu^a | a \rangle = 0. \quad (14)$$

Now we apply these to our study. Let’s the PQ charges of the relevant scalar fields $h, h', S_1, S_2$ be $\Gamma^{(h)}, \Gamma^{(h')}, \Gamma^{(1)}, \Gamma^{(2)}$, respectively. The PQ current is

$$j_\mu^a = i(\sum_i \Gamma_i^{(h)} h_i^u \partial_\mu h_i^u + \sum_j \Gamma_j^{(h')} h_j^d \partial_\mu h_j^d + \sum_k \Gamma_k^{(1)} S_{1k}^u \partial_\mu S_{1k}^u + \sum_l \Gamma_l^{(2)} S_{2l}^u \partial_\mu S_{2l}^u) + (\text{extra singlet and quark/lepton currents})
$$

$$= \sum_i \Gamma_i^{(h)} v_i^u \partial_\mu \xi_i^u + \sum_j \Gamma_j^{(h')} v_j^d \partial_\mu \xi_j^d + \sum_k \Gamma_k^{(1)} u_k^{(1)} \partial_\mu \xi_k^{(1)} + \sum_l \Gamma_l^{(2)} u_l^{(2)} \partial_\mu \xi_l^{(2)} + (\text{extra singlet and quark/lepton currents})
$$

$$= f_a \partial_\mu a + (\text{extra singlet and quark/lepton currents}), \quad (15)$$

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where \( i, j, k, l \) are the generation indices. \( v^{(u)}, v^{(d)}, u^{(1)} \) and \( u^{(2)} \) are the VEVs of \( h, h', S_1 \) and \( S_2 \), respectively. If there are some extra singlets, those contributions should be included as indicated in the parentheses. The axion field is

\[
a = \frac{1}{f_a} \left( \sum_i \Gamma_i^{(h)} v_i^{(u)j} \xi_i^{(u)} + \sum_j \Gamma_j^{(h')} v_j^{(d)j} \xi_j^{(d)} + \sum_k \Gamma_k^{(1)} u_k^{(1)k} \xi_k^{(1)} + \sum_l \Gamma_l^{(2)} u_l^{(2)l} \xi_l^{(2)} \right),
\]

where

\[
f_a^2 = \sum_i (\Gamma_i^{(h)} v_i^{(u)})^2 + \sum_j (\Gamma_j^{(h')} v_j^{(d)})^2 + \sum_k (\Gamma_k^{(1)} u_k^{(1)})^2 + \sum_l (\Gamma_l^{(2)} u_l^{(2)})^2.
\]

We can generally consider the gauge structure \( SU(3)_C \times SU(2)_L \times U(1)_Y \times U(1)^2(\subset E_6) \) and also take the \( \chi, \psi \) basis with respect to extra \( U(1)^2 \) as defined in (8). There are three neutral currents \( j_\mu^Z, j_\mu^\chi, j_\mu^\psi \) which couple to \( Z^0 \) and extra \( U(1)^2 \) gauge bosons, respectively. The axion should decouple from these currents,

\[
\langle 0 | j_\mu^{Z^0} | a \rangle = \langle 0 | j_\mu^\chi | a \rangle = \langle 0 | j_\mu^\psi | a \rangle = 0.
\]

The axion should also be orthogonal to \( \pi_0 \). But here we donot discuss this condition. If there are more extra gauges, the additional decoupling conditions must be imposed on the axion field \( a \). The neutral currents are expressed as

\[
j_\mu^{Z^0} = \frac{1}{v} \left( \sum_i v_i^{(u)i} \partial_\mu \xi_i^{(u)} - \sum_j v_j^{(d)j} \partial_\mu \xi_j^{(d)} \right),
\]

\[
j_\mu^\chi = \frac{1}{w} \left( \frac{2}{\sqrt{6}} \sum_i v_i^{(u)i} \partial_\mu \xi_i^{(u)} - \frac{2}{\sqrt{6}} \sum_j v_j^{(d)j} \partial_\mu \xi_j^{(d)} - \frac{5}{\sqrt{6}} \sum_k u_k^{(1)k} \partial_\mu \xi_k^{(1)} \right),
\]

\[
j_\mu^\psi = \frac{1}{w} \left( \frac{-2}{\sqrt{3}} \sum_i v_i^{(u)i} \partial_\mu \xi_i^{(u)} - \frac{2}{\sqrt{3}} \sum_j v_j^{(d)j} \partial_\mu \xi_j^{(d)} + \frac{1}{\sqrt{3}} \sum_k u_k^{(1)k} \partial_\mu \xi_k^{(1)} + \frac{4}{\sqrt{3}} \sum_l u_l^{(2)l} \partial_\mu \xi_l^{(2)} \right).
\]

\( u, v \) and \( w \) are defined in the similar way as eq.(13). Substituting eqs.(16) and (19) into eq.(18), we get the decoupling conditions

\[
\sum_i \Gamma_i^{(h)} v_i^{(u)^2} - \sum_j \Gamma_j^{(h')} v_j^{(d)^2} = 0,
\]

\[
\sum_k \Gamma_k^{(1)} u_k^{(1)^2} = 0,
\]

\[
\sum_l \Gamma_l^{(2)} u_l^{(2)^2} = 0.
\]

\( ^6 \)It should be noted that there are necessarily two extra \( U(1) \) factors in the models with the intermediate mass scale \( [1] \) whose existence is the necessary condition for our mechanism.
In our scenario the PQ symmetry is built in the theory from the beginning. Therefore there is no freedom to tune the PQ charge so as to guarantee the existence of the axion decoupling from the neutral gauge currents correctly. Depending on the model it is determined automatically whether the invisible axion exists or not.

To study this problem, we need to determine the PQ charge of each scalar fields $\phi_i$. From the fact that the phase part of eq.(2) corresponds to $U(1)_{PQ}$, PQ charge of $\phi_i$ is

$$\Gamma_i = n_i \tan^{-1}\left(\frac{T_R}{\alpha - T_I}\right).$$

(21)

$\alpha$ is an arbitrary parameter and $T_R, T_I$ are real and imaginary parts of $T$, respectively. $\tan^{-1}\left(\frac{T_R}{\alpha - T_I}\right)$ is considered as the normalization factor of the PQ charge. For the untwisted matter fields $n_i$ equals to $-1$ and the twisted matter fields have $n_i \leq -1$ integer values depending on the way of the twists. Calabi-Yau type models have no massless twisted matter fields. Therefore the modular weights have the same sign and then the PQ charges have also the same sign for all matter fields. As a result the axion decoupling condition (20) cannot be satisfied realizing the non-trivial hierarchy

$$v^{(u)} \sim v^{(d)} < u^{(1)} \ll u^{(2)}.$$  

(22)

This hierarchical structure is necessary to bring the symmetry breaking pattern of the standard model\cite{12}. Thus it is difficult for our axion mechanism to work in the Calabi-Yau type models.

Usually in the 4-dimentional superstring there exist extra $U(1)$ factors and new matter singlets other than those of the $E_6$ models and we can expect theses ingredients to remedy the situation discussed above in the Calabi-Yau case. In that case the restriction on the superpotential due to the discrete symmetry will be necessary. For example, the role of extra gauge symmetry in (ii) should be played by the discrete symmetry.

Next we study the compatibility of the proton stability and the axion cosmological bound. As shown in eq.(9), $g$ and $\bar{g}$ couple to the ordinary quarks and leptons. And these couplings induce the proton decay through the tree level couplings $O(\frac{\lambda_4 \lambda_5}{M_g^2})(QQQ\ell)$ and/or $O(\frac{\lambda_5 \lambda_6}{M_g^2})(\bar{u}\bar{u}\bar{d}\bar{e})$\cite{12} \cite{13}. Since Yukawa coupling constants $\lambda_i$ are usually order one\cite{14}, these amplitudes are determined by $M_g$ which is estimated through $S_2\bar{g}g$ coupling as $M_g = \ldots$

\footnote{We should note that $S_1$ cannot have a VEV at high energy scale like $S_2$ because $S_1$ has no D-term flat potential. See also the arguments in the next paragraph.}

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Here the intermediate mass scale $\langle S_2 \rangle$ is introduced in the following way. The scalar potential of $S_2$ is composed of F- and D-terms. In order to keep the supersymmetry until the weak scale, $S_2$ must have the flat potential at the intermediate scale. D-term flatness is guaranteed by $\langle S_2 \rangle = \langle \bar{S}_2 \rangle$ because the D-term contribution to the scalar potential is $\sum_\alpha (S_2^\dagger T_\alpha S_2 - \bar{S}_2^\dagger T_\alpha \bar{S}_2)$ where $T_\alpha$ expresses the $U(1)$ charge. Taking account of the absence of $S_3^2$ term in the superpotential $W$ due to the extra $U(1)$ symmetry, the scalar potential are produced by the non-renormalizable terms

$$V = \lambda(p) M_C^{6-4p} S_2^{4p-2} - M_S^2 S_2^2,$$  \hspace{1cm} (23)

where $M_C$ is the compactification scale and $p$ represents the lowest order contribution. $M_S^2 S_2^2$ is the soft supersymmetry breaking term induced from the hidden sector. From this we get

$$\langle S_2 \rangle = M_C \left( \frac{M_S}{M_C} \right)^{(\alpha-1)/\alpha}.$$  \hspace{1cm} (24)

The present experimental bound of the proton stability requires $\langle S_2 \rangle > 10^{16}\text{GeV}$. This is realized only when $p \geq 4$. And such a condition is known to be satisfied in some models which have the large symmetry. The cosmological bound on the axion model imposes on the axion decay constant $f_a < 10^{12}\text{GeV}$. If $f_a \sim \langle S_2 \rangle$ as expected from eq.(13), the discrepancy occurs between the proton stability and the axion cosmological bound.

In order to make our axion mechanism realistic we must reconcile these. The most simple solution for this will be the following one. The tree level amplitude of the proton decay is zero because the relevant Yukawa couplings in the superpotential $W$ are zero. And $f_a \sim \langle S_2 \rangle \sim \sqrt{M_C M_S} \sim 10^{11}\text{GeV}$ ($p = 2$). This possibility has been studied for the proton stability in the various superstring models. But within our knowledge the realistic model which realizes this has not been found by now.

We may consider the other type solution for this problem. As mentioned before, the axion decay constant $f_a$ is not the measure of the symmetry breaking scale. They are related by eq.(13) and in our case PQ charge is given by eq.(21). From the study of the effective theory with target space modular invariance, we know $\langle T_I \rangle = 0$, $\langle T_K \rangle = O(1)$ and then $\Gamma_i \sim \eta_i \tan^{-1}(1/\alpha)$. If this PQ charge normalization is taken to be extremely small (i.e. $\alpha \sim 10^5$), we can realize simultaneously $f_a < 10^{12}\text{GeV}$ and $M_g > 10^{16}\text{GeV}$ in the model with $\langle S_2 \rangle > 10^{16}\text{GeV}$. However, this appears to be unnatural. The reconciliation
of these may remain potentially as a serious problem for the present mechanism not only in Calabi-Yau type models but also other 4-dimensional superstring models unless we find the way to suppress the coupling of the extra colored fields to ordinary quarks and leptons.

In the above consideration we assume the existence of the extra color triplets $g$ and $\bar{g}$ as suggested by Ibáñez and Lüst. However, following eq.(10) DFSZ type invisible axion mechanism seems to be able to work without $g$ and $\bar{g}$. If there are no $g$ and $\bar{g}$, the proton stability problem related to the intermediate mass scale $\langle S_2 \rangle$ will disappear and we can take $\langle S_2 \rangle \sim 10^{11}\text{GeV}$ which is consistent with the cosmological bound. Moreover there may be a good feature that the so-called $\mu$-problem[17] does not exist because there are generally two kinds of singlet fields like $S_2$ as stressed in ref.[12]. Only $S_2$ which has the partner $\bar{S}_2$ can have a VEV at the intermediate scale because of the D-term flatness. $S_2$ which has no partner $\bar{S}_2$ remains massless until the weak scale and it can contribute to $\lambda_9$ term which is relevant to the symmetry breaking at the weak scale. In the viewpoint of proton stability this possibility may be more realistic than KSVZ type solution in which the extra colored heavy fermions play the important role. Unfortunately in Calabi-Yau type models we cannot find such a solution because $g$ and $\bar{g}$ are contained in 27 of $E_6$ and remain massless at the compactification scale. Orbifold may give such models.

In conclusion, we applied the solution for the strong CP problem in the 4-dimensional superstring theory recently proposed by Ibáñez and Lüst to Calabi-Yau type models. Their mechanism is very interesting but in the Calabi-Yau type models there seem to be difficult phenomenological problems. It is very interesting to find an explicit $(0, 2)$ orbifold models in which their mechanism can be realized in phenomenologically successful way. The DFSZ type invisible axion model which works without heavy extra colored $g$ and $\bar{g}$ fields will be more promising than KSVZ axion from the viewpoint of the proton stability. It is also interesting subject to construct such models concretely.
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