Low-Scale Anomalous U(1) and Decoupling Solution to Supersymmetric Flavor Problem

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Abstract

Supersymmetric standard models where the ultraviolet cut-off scale is only a few orders of magnitude higher than the electroweak scale are considered. Phenomenological consequences of this class of models are expected to be very different from, for example, the conventional supergravity scenario. We apply this idea to a model with an anomalous U(1) gauge group and construct a viable model in which some difficulties of the decoupling solution to the supersymmetric flavor problem are ameliorated.

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What is the fundamental scale of the unified theory of particle physics? Usually it is supposed to be near the Planck scale. In fact this view is supported by the weakly coupled heterotic string theory where the string scale must be close to the Planck scale \([1]\). It has recently been recognized, however, that the fundamental scale can be much lower than the Planck scale, even as low as around 1 TeV \([2]\). In this case, the largeness of the Planck scale is accounted for by extra dimensions with large volume, in which gravity propagates, while the standard model particles have to be confined on four dimensional subspace (3-brane). Remarkably this configuration survives various phenomenological constraints \([2–5]\). Furthermore it can be realized in, for example, type I (or more precisely type I’) superstring theory where the string scale may not directly be related to the Planck scale \([3]\) (See also Ref. \([7]\) for earlier attempts). Even when the fundamental scale itself is close to the Planck scale, special geometrical configuration of extra dimension(s) such as an AdS slice \([8]\) may enable us to obtain a model with effectively very low cut-off scale on a visible brane where the standard model sector is confined.

In this paper, we would like to consider the situation where the fundamental scale, or the ultraviolet cut-off scale, is much lower than the Planck scale but still lies a few order of magnitude above the electroweak scale so that low energy supersymmetry is needed to protect the electroweak scale from radiative corrections. A supersymmetric standard model with such a low-scale cut-off will naturally fall into a class of low-scale supersymmetry breaking models. An immediate consequence of this class of models is that the gravitino, the superpartner of the graviton, is much lighter than other superparticles, and thus tends to be the lightest superparticle. Moreover soft supersymmetry breaking masses are given at the low energy scale so that the mass spectrum is in general quite different from that of high-scale supersymmetry breaking models.

Inspired by the arguments given above, here we would like to discuss phenomenological implications of supersymmetric standard models with such a low fundamental scale. Specifically we consider a model with anomalous \(U(1)_X\) gauge symmetry \([9–11]\). It is well-known that appropriate assignment of the \(U(1)_X\) charges provides the decoupling solution of the su-
persymmetric flavor problem \cite{12}. Namely the first two generations of squarks and sleptons are assumed to be heavy, thus suppress flavor changing neutral current (FCNC) processes from superparticle loops, while squarks and sleptons in the third generation are set to be at the electroweak scale, \textit{i.e.} a few hundred GeV: otherwise the large Yukawa couplings for the third generation would generate large radiative corrections to the Higgs mass and then we would lose the very motivation to introduce low-energy supersymmetry. The decoupling solution is simple and attractive, in particular when symmetries relate the smallness of the masses of the quarks and leptons in the first two generations with the largeness of the masses of the squarks and sleptons in the first two generations.

However it has been pointed out that there are several difficulties in this scenario. First, although the squarks and sleptons in the first two generations do not influence the running of the Higgs mass at one-loop level if one ignores small Yukawa couplings, they do at two-loop level. Thus they cannot be arbitrarily heavy: rather their masses are severely limited by the naturalness argument. In fact, this issue was discussed in detail in Ref. \cite{13} which gave an upper bound of 5 TeV from the condition that the fine tuning should be less than 10\%. The second problem is that the heavy squarks would give negative contribution to the mass squared of the third generation squarks at two-loop level, driving the mass squared negative to cause color breaking \cite{14,15}. It turns out that the bound is very severe. In fact the masses of the sfermions in the first two generations must be much smaller than 10 TeV and thus they are not large enough to solve the FCNC problem, as far as the mass of the stop is lighter than about 1 TeV. Finally with a mass spectrum typical in the decoupling solution, the relic abundance of the lightest superparticle which is assumed to be bino-like neutralino tends to overclose the Universe \cite{16}.

What we will do in this paper is first to construct a viable model with the low fundamental scale and then to show that the problems above are ameliorated in such a framework. Here we should note that the issue of the color instability in a scenario where supersymmetry breaking is mediated at a substantially low energy was discussed in Ref. \cite{15}, but without an explicit model. Also a different approach to the problem of the color instability by adding
extra matter multiplets to eliminate dangerous contributions has been proposed in Ref. [17].

We begin by describing the model we are considering. The model is similar to that of Dvali and Pomarol [9] (See also Ref. [18]). Ref. [9] considered a high scale cut-off theory such as the heterotic string, where the Fayet-Iliopoulos (FI) term for the anomalous $U(1)_X$ gauge group is given by $\xi = \frac{\pi Q g^2}{16 \pi^2} M_P^2$ with $g$ the gauge coupling constant. On the contrary, what we will consider is a low-scale cut-off theory with the cut-off $M_*$. Here we assume that the standard model sector as well as the $U(1)_X$ is confined on a brane-like object such as a D-brane and the large four dimensional Planck scale requires the existence of large extra dimensions in which gravity propagates. Then it is natural to expect that the FI term $\xi$ is, if it is non-zero,

$$|\xi| \leq M_*^2. \quad (1)$$

From now on we will take a convention $\xi > 0$. $M_*$ may be identified with the string scale. In type I and type IIB string models, $\xi$ will be generated through non-vanishing expectation values of some moduli fields when combined with a generalized Green-Schwarz anomaly cancellation mechanism [19,20].

As for chiral multiplets, we introduce $\phi_+$ and $\phi_-$ with $U(1)_X$ charge +1, and −1, respectively, and $y_i$ with charge $Q_i$, which represent fields in the standard model sector. Then the $U(1)_X$ D-term is written

$$D = \xi + |\phi_+|^2 - |\phi_-|^2 + \sum_i Q_i |y_i|^2 \quad (2)$$

The model also has the following mass term in the superpotential

$$W = m \phi_+ \phi_- \quad (3)$$

besides the superpotential of the standard model sector. Here we assume $m^2 \leq g^2 \xi$. Though it is possible to generate the mass term Eq.(3) dynamically, we will treat it as a given parameter. Note that this does not mean to introduce a huge hierarchy into mass parameters, since all the mass scales of this low-scale theory are not very far from the electroweak scale.
By minimizing the scalar potential of the model

\[ V = \left| \frac{\partial W}{\partial \phi^+} \right|^2 + \left| \frac{\partial W}{\partial \phi^-} \right|^2 + \frac{g^2}{2} D^2 \]  

(4)

with \( g \) being the gauge coupling constant for \( U(1)_X \), we find the following vacuum expectation values

\[
\phi^+ = 0, \\
\phi^- = \sqrt{\xi - \frac{m^2}{g^2}}, \\
F_{\phi^+} = m \phi^- = m \sqrt{\xi - \frac{m^2}{g^2}}, \\
F_{\phi^-} = 0, \\
D = \frac{m^2}{g^2}
\]

(5)

Here we have neglected the contributions from the standard model sector, which are assumed to be tiny.

In this model the scalar masses in the standard model sector are written in the following form:

\[ m_0^2 = Q g^2 D + m_F^2 = Q m^2 + m_F^2, \]

(6)

which are given at the cut-off scale \( M_\star \). The first term is the \( U(1)_X \) D-term contribution which is solely controlled by the \( U(1)_X \) charge \( Q \). On the other hand, the second term which represents a F-term contribution comes from non-renormalizable interaction in the Kähler potential and is sensitive to the physics close to the cut-off scale. In fact we expect to have the following term in the Kähler potential:

\[ \frac{\eta_{ij}}{M_\star^2} \phi^+ \phi^- y_i^* y_j \]

(7)

with numerical coefficients \( \eta_{ij} \) of order unity or less, which are in general generation dependent.\(^1\) Eq. (7) yields a (possibly) generation dependent F-term mass estimated at most

\(^1\)Non-renormalizable interactions including bulk fields will be suppressed by the four dimensional Planck mass \( M_{Pl} \).
\[ m_{P}^{2} \sim \frac{F_{\delta_{+}}^{2}}{M_{*}^{2}} \sim m^{2}\left(\frac{\xi - m^{2}/g^{2}}{M_{*}^{2}}\right). \]  

(8)

Therefore this is potentially a source for the FCNC. However it is always sub-dominant compared to the first term, provided that there is a little hierarchy of one order of magnitude or so between \(\sqrt{\xi - m^{2}/g^{2}}\) and \(M_{*}\):

\[ \epsilon \equiv \frac{\sqrt{\xi - m^{2}/g^{2}}}{M_{*}} \leq O(10^{-1}). \]  

(9)

Here it is interesting to mention the mass spectrum of \(\phi_{+}\) and \(\phi_{-}\). What happens is that both supersymmetry and the \(U(1)_{X}\) gauge symmetry are broken spontaneously. Since the scalar component of \(\phi_{-}\) acquires non-zero vacuum expectation value, its real component has a similar mass to the gauge boson mass, \(\sqrt{g^{2}\xi - m^{2}}\), and its imaginary component becomes the would-be Nambu-Goldstone boson. On the other hand, it is essentially the \(\phi_{+}\) multiplet which is responsible for supersymmetry breaking, and its spinor component is the Goldstino absorbed into the gravitino in supergravity framework. In this case the mass of the scalar component of the \(\phi_{+}\) multiplet is found to be \(\sqrt{2m}\), similar in size with the soft supersymmetry breaking masses.

Let us now discuss how the low cut-off scale model ameliorates the problems of the decoupling solution. First we will consider the color instability. For simplicity, we assign \(U(1)_{X}\) charges for quarks and leptons in the first-two generations to be +1 and those for other matters to be 0. In this framework, from the cut-off scale \(M_{*}\) to the mass scale of the sfermions in the first two generations \(\tilde{m}_{1,2}\) the nature can be described by four-dimensional field theory with the matter content of the minimal supersymmetric standard model (MSSM), and from the \(\tilde{m}_{1,2}\) scale to about 1 TeV it can be described by an effective theory in which squarks and sleptons in the first-two generations are integrated out.

A constraint on the soft masses in the third generation at the cut-off scale is obtained by requiring physical masses to be positive at the electroweak scale. Although the physical masses receive D-term contributions of the \(SU(2)_{L} \times U(1)_{Y}\) gauge interactions and also
left-right mixing effects, we neglect them for simplicity (the effects due to these neglect are discussed in [15]). Hereafter we will obtain a constraint by requiring the running masses to be positive at 1 TeV scale as was done in [14,15].

The values of soft masses at 1 TeV scale are computed by using renormalization group equations (RGEs). In our analysis, we use the two-loop RGEs in the $\overline{\text{DR}}$ scheme [21]. For our charge assignment, the RGEs for the soft masses in the third generation that include Yukawa couplings, $A$-terms at one-loop level and heavy sfermion contributions at two-loop level in this scheme are

$$
\mu \frac{d\tilde{m}_f^2}{d\mu} = -\frac{2}{\pi} \sum_A \alpha_A C_A^f M_A^2 + \frac{1}{2\pi} \eta_f^f \alpha_t (\tilde{m}_{Q_3} + \tilde{m}_{U_3} + \tilde{m}_{H_u} + A_t^2)
$$

$$
+ \frac{1}{2\pi} \eta_f^b \alpha_b (\tilde{m}_{Q_3} + \tilde{m}_{D_3} + \tilde{m}_{H_d} + A_b^2) + \frac{1}{2\pi} \eta_f^\tau \alpha_b (\tilde{m}_{L_3} + \tilde{m}_{E_3} + \tilde{m}_{H_d} + A_\tau^2)
$$

$$
+ \frac{2}{\pi^2} \sum_A \alpha_A^2 C_A^f \tilde{m}_{1,2}^2,
$$

where $\alpha_A$ and $C_A^f$ are the gauge couplings and the quadratic Casimir of $SU(3)_C$, $SU(2)_L$ and $U(1)_Y$ for $A = 3, 2, 1$, respectively, and $\alpha_f = Y_f^2/4\pi (f = t, b, \tau)$ are Yukawa couplings, with $\eta_f^f = (1, 2, 3)$ for $f = Q_3, U_3, H_u$, $\eta_f^b = (1, 2, 3)$ for $f = Q_3, D_3, H_d$, and $\eta_f^\tau = (1, 2, 1)$ for $f = L_3, E_3, H_d$, respectively, and zero otherwise. Note that the contribution proportional to the $U(1)_Y$ D-term does not appear in (10) for the sake of the boundary conditions for the squark and slepton masses. We solve the RGEs as follows. At the cut-off scale, we assume that gaugino masses satisfy

$$
M_3(M_*) = \frac{\alpha_3(M_*)}{\alpha_2(M_*)} M_2(M_*) = \frac{\alpha_3(M_*)}{\alpha_1(M_*)} M_1(M_*)
$$

for simplicity. Below the $\tilde{m}_{1,2}$ scale, the heavy sfermions do not contribute to the running of the couplings and masses. Note that the gaugino masses evolve differently from the gauge coupling constants, but we ignored this deviation which is not important in our analysis.

For the squark sector the gluino contribution dominate the other gaugino contributions, so that the constraint for the squark masses is insensitive to this assumption. Because of the charge assignment mentioned above, the squarks and sleptons in the first-two generations have soft masses $\tilde{m}_{1,2} \simeq m$ at the cut-off scale. The soft masses of the third generation
scalar bosons and the Higgs bosons have no contribution from the $U(1)_X$ D-term and their F-term contributions are model dependent. Here to simplify the analysis we assume them all universal: $m_0 \simeq \epsilon \tilde{m}_{1,2} \simeq \epsilon m$ at the cut-off scale. We take into account the bottom and tau Yukawa couplings as well as the top Yukawa coupling. In our analysis, we fix the top quark mass in the $\overline{\text{MS}}$ scheme $m_t^{\overline{\text{MS}}}(m_t)$ to be 167 GeV and that of the bottom quark $m_b^{\overline{\text{MS}}}(m_b)$ to be 4.3 GeV. We checked whether our results depend on $\tan \beta$ and $A$ parameters and found that the dependence on these parameters are very small. Thus we will present results with all $A$ parameters zero at the cut-off scale and $\tan \beta = 2$.

We also have to include finite term contribution because the scale at which the initial condition of RGEs is given is not much larger than the electroweak scale. We follow Ref. [22] to evaluate this effect and the result for our case is

$$\tilde{m}_{f,\text{finite}}(\mu) = -\frac{1}{\pi^2} \left( \frac{\pi^2}{3} - 2 - \ln \left( \frac{\tilde{m}_{1,2}}{\mu^2} \right) \right) \sum_A \alpha^2_A C_A f \tilde{m}_{1,2}^2. \quad (12)$$

At the $\tilde{m}_{1,2}$ scale, we add this contribution to the running mass, as a threshold effect. The $U(1)_Y$ D-term contribution is absent as in the case of Eq. (10). Note that the finite contribution (12) is different from that of Ref. [15]. The difference can be absorbed by a redefinition of the renormalization scale $\mu$.

In fig. 1, we show the allowed maximum values of the sfermion masses in the first two generations $\tilde{m}_{1,2}$ by requiring that the mass squared $\tilde{m}_{Q_3}^2$ be positive at the 1 TeV scale. The horizontal axis represents the running gluino mass at the scale $\mu = 1$ TeV. Here we take all the scalar masses of the third generation and the soft masses of the Higgs doublets at the cut-off scale $\tilde{m}_f(M_*)$ to be $1$ TeV. We found that the constraints from the positivity requirements of $\tilde{m}_{L_3}^2$ and $\tilde{m}_{D_3}^2$ are similar to that from $\tilde{m}_{Q_3}^2$ presented here, and in fact the differences are less than 10%. We consider the cases $M_* = 10^3$, $10^4$ and $10^5$ TeV. For comparison, we also show the case $M_* = 10^{16}$ GeV. The region above each curve is excluded. As $M_*$ decreases, one finds that the allowed region becomes larger because the RGE effect becomes less significant. Indeed the finite term dominates when $M_* = 100$ TeV. When the gluino mass is, for instance, about 1 TeV, $\tilde{m}_{1,2}$ can be as heavy as about 17 TeV for
\( M_* = 100 \, \text{TeV} \). In this case the contribution to the kaon mass difference \( \Delta m_K \) will become at an acceptable level with the help of a small alignment between squark mass eigenstates and interaction eigenstates \[23\]. The constraint from the positivity requirement of the third generation slepton \( \tilde{m}_{L_3}^2 \) alone is not so strong. For example, \( \tilde{m}_{1,2} \) is required to be smaller than about 80 TeV and 25 TeV, for \( M_* = 100 \, \text{TeV} \) and \( M_* = 10^5 \, \text{TeV} \), respectively, as far as the gluino mass at the 1 TeV is smaller than 3 TeV. On the other hand, in the case \( M_* = 10^{16} \, \text{GeV} \) \( \tilde{m}_{1,2} \) is required to be smaller than about 13 TeV.

Next we would like to discuss how the other difficulties are cured in our setting. The point of the naturalness problem discussed in Ref. \[13\] is that the heavy first two generation scalar masses will influence the running of the Higgs mass at two loop level, causing the fine tuning to obtain the electroweak scale if the masses are very heavy. Now since the contribution to the running is roughly proportional to the "length" of running in logarithmic scale, the fine-tuning problem should be relaxed in our low-scale cut-off case in which the length of running is much shorter than the high-scale cut-off case discussed by \[13\].

In our scenario, the gravitino becomes very light with the estimate

\[
m_{3/2} = \frac{F_{\phi^+}}{\sqrt{3} M_{Pl}} \simeq m \frac{\sqrt{\xi - m^2/g^2}}{M_{Pl}} \simeq 0.1 \text{eV} \left( \frac{M_*}{100 \, \text{TeV}} \right) \left( \frac{\epsilon m}{1 \, \text{TeV}} \right).
\]

Thus it is likely to be the lightest superparticle (LSP). Then the lightest superpartner in the supersymmetric standard model sector is no longer stable, but it immediately decays into the gravitino and hence it is obvious that the overclosure problem of the neutralinos is evaded. In this scenario, the gravitino is stable. Its cosmological implications are discussed in the literature. For the gravitino which weighs much less than 1 keV, its relic abundance is much smaller than the critical density of the Universe and thus it is cosmologically harmless. See Ref. \[24\] and references therein for detail. We should also note that superparticle signals at collider experiments in our scenario have some characteristic features. The lifetime of the next to the lightest superparticle (NSP) is roughly of the order

\[
\tau_{\text{NSP}} \simeq 16\pi \frac{F_{\phi^+}^2}{m_{\text{NSP}}^5} \simeq 10^{-17} \text{sec} \left( \frac{100 \, \text{GeV}}{m_{\text{NSP}}} \right)^5 \left( \frac{\epsilon m}{1 \, \text{TeV}} \right)^2 \left( \frac{M_*}{100 \, \text{TeV}} \right)^2,
\]

(14)
assuming that the decay is a two-body decay. Thus the lifetime is so short that it will decay inside a detector. If the NSP is bino-like, the decay contains a photon and a gravitino which escapes detection. If the NSP is a slepton, most likely a stau, the decay contains a tau lepton and a gravitino. In our scenario, the stop may be the NSP. In this case, the stop decays into a top (or a W boson and a bottom quark\footnote{For the three-body decay, the decay length increases substantially. An analysis in this case has been given in \cite{25}.}) and a gravitino. In either case the signals will be distinguishable from those of high-scale supersymmetry breaking scenario where the signals will be associated with a massive LSP escaped from a detector. Here it should also be noticed that we can probe the heavy mass scale of the first-two generations via superoblique corrections \cite{26}, even though they cannot be produced directly in near future colliders.

Here we would like to briefly mention a gaugino mass. In our model we can write the following term

\[
\frac{\phi^+ \phi^-}{M^2_s} W^\alpha W_\alpha, \tag{15}
\]

where \( W^\alpha \) is a supersymmetric field strength of a gauge field. It follows from this that the gaugino mass is of the order

\[
\frac{F_{\phi^+ \phi^-}}{M^2_s} = \epsilon^2 m. \tag{16}
\]

Recall that the mass of the third generation squark is \( \sim \epsilon m \). An additional suppression factor \( \epsilon \) in Eq. (16) may be compensated by unknown numerical coefficients in front. Note that in the low-scale supersymmetry breaking scenario, the contribution to the gaugino mass from superconformal anomaly \cite{27} is negligible because it is proportional to the tiny gravitino mass.

Before concluding we will comment on the large extra dimensions needed to obtain the large Planck scale with the low string scale scenario. The size of the compact \( n \)-dimensional extra dimensions \( R \) is given by
\[ M_{Pl}^2 \simeq M_*^{2+n} R^n, \quad (17) \]

or

\[ R^{-1} \simeq M_* \left( \frac{M_*}{M_{Pl}} \right)^{2/n} \quad (18) \]

to reproduce the Planck scale \( M_{Pl} \). Here we have assumed that the compact manifold is isotropic and is characterized by a single size \( R \). To illustrate, let us take \( n = 6 \). Then

\[ R^{-1} \simeq 10^{2} \text{GeV} \left( \frac{M_*}{10^6 \text{GeV}} \right)^{4/3}. \quad (19) \]

The masses of graviton’s Kaluza-Klein modes are quantized in units of \( R^{-1} \). Thus we find that the KK mode masses are in the electroweak scale or higher. This contrasts with the case of the large extra dimension scenario with \( M_* \simeq 1 \text{ TeV} \) where \( R^{-1} \simeq 10 \text{ MeV} \).

Since the KK modes have masses of the electroweak scale or so and have interactions similar to the graviton, they may affect cosmological evolution of the early Universe. In particular, they are produced after the inflationary epoch and decay typically around the epoch of the big-bang nucleosynthesis. Here we will not go into detailed discussion, but make some remarks. First if \( R^{-1} > 10^4 \text{ GeV} \), the KK modes decay before the nucleosynthesis and thus they are harmless. On the other hand, for \( 10^2 \text{ GeV} < R^{-1} < 10^4 \text{ GeV} \), the reheat temperature after inflation must be low to suppress the production of the KK modes. We expect that the reheat temperature of \( 10^2 \text{ GeV} \) will be allowed since then the production of the KK modes whose masses are heavier than \( 10^2 \text{ GeV} \) is highly suppressed. This is very different from the TeV gravity case where the reheat temperature is forced to be even smaller than 1 GeV [2]. The higher reheat temperature in our case has an advantage for baryogenesis. In particular one may be able to use the electroweak baryogenesis.

It is interesting to mention here that the radius of the extra dimension can be as large as a sub-millimeter for \( n = 1 \) and \( M_* \simeq 10^8 \text{ GeV} \). This case may be tested in a future gravity experiment [2].

There remains a problem of how one realizes a viable inflation model and a subsequent stabilization of the size of the extra dimensions. A hope is that model building for this may
be somewhat easier than the original large extra dimension scenario [28]. This issue should deserve further study.

To summarize, we have considered the supersymmetric standard model with the anomalous $U(1)$ gauge symmetry when the ultraviolet cut-off scale is not far from the electroweak scale. In our scenario, the Fayet-Illiopoulos D-term is set to be a bit smaller than the cut-off scale squared. Except this, the model is similar to that of [1]. We applied this model to the decoupling solution of the supersymmetric flavor problem and showed that the difficulties of the solution become less severe than the conventional high-scale cut-off scenario. The model should be combined with the idea of the large extra dimensions to obtain the large Planck scale of the gravitational interaction. We briefly discussed some of the related cosmological issues.

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FIG. 1. The allowed maximum values of the sfermion masses in the first two generations $\tilde{m}_{1,2}$ by requiring that the stop mass squared $\tilde{m}_{Q_3}^2$ be positive at the 1 TeV scale, for the cut-off scale $M_*= 100, 10^3, 10^4, 10^5$ TeV and $M_* = M_{\text{GUT}} = 10^{16}$ GeV. The region above each curve is excluded. The horizontal axis represents the running gluino mass at the scale 1 TeV. All the scalar masses of the third generation as well as the soft masses of the Higgs doublets at the cut-off scale are fixed to be 1 TeV.