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Model-dependent axion as quintessence with almost massless hidden sector quarks

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Abstract: A pseudo-Goldstone boson for *quintessence* is known to require the decay constant $F_q \sim M_P$ and the height of the potential $(0.003\,\mathrm{eV})^4$, and hence the pseudo-Goldstone boson mass of order 10^{-33} eV. The model-dependent axion in some superstring models is suggested as a good candidate for the quintessence. We realize this idea with a hidden-sector gauge group $\mathrm{SU}(3)_H \times \mathrm{U}(1)_H$. The hidden-sector instanton contribution to the quintessence potential is suppressed by the almost massless hidden-sector quarks. The model-independent axion is the QCD axion required for the strong CP solution.

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Various mass scales in physics have been the most unambiguous hints toward new physics. For example, the weak scale $\sim 250\,\mathrm{GeV}$ has led to the construction of the standard model. The Einstein-Hilbert action with the scale parameter $M_P = 2.44 \times 10^{18}\,\mathrm{GeV}$ gives the long range gravitational interaction. From the huge ratio of these two known scales, technicolor and supersymmetry solutions of the hierarchy problem were proposed [1], which triggered the main research activities of high energy physics in the last twenty years. In addition there exists a theoretically favored (but still experimentally unproven) axion scale at the window $10^{9-12}\,\mathrm{GeV}$ [2] which awaits to be uncovered in the near future [3]. Also, there exists the scale μ in supersymmetric models which is in principle an independent parameter [4].

In addition, another very important scale revealed gradually in the recent years from the shadow is the vacuum energy density or the cosmological constant. Since the Einstein gravity without the cosmological term was so successful, it has been the theoretical prejudice that it must vanish. Experimentally, one could have given at most an upper bound on the cosmological constant before 1998. But with the first round observation on the negative decceleration parameter in 1998 [5], which has to be confirmed with more data, the hypothesis of a positive but a small non-zero cosmological constant has attracted a great deal of attention [6]. If confirmed, it introduces a new physical scale of order

$$\Lambda_{c.c.} = (0.003\,\mathrm{eV})^4 \simeq 10^{-46}\,\mathrm{GeV}^4\,. \tag{1}$$

This scale for the cosmological constant is so small that it must come from a ratio of a hierarchical number or from an exponentially suppressed number. From the perturbative analysis, one usually obtains a hierarchical ratio. If so, a plausible ratio is

$$\frac{v_{\rm ew}^7}{M_P^3} \sim 4.2 \times 10^{-39}\,\mathrm{GeV}^4\,, \qquad \frac{v_{\rm ew}^8}{M_P^4} \sim 4.3 \times 10^{-55}\,\mathrm{GeV}^4\,, \tag{2}$$

where $v_{\rm ew}$ is the known electroweak scale. The term suppressed by M_P^4 gives a too small number responsible for the non-zero cosmological constant since the possible couplings would reduce the number further. The term suppressed by M_P^3 is a bit large but the couplings multiplied would lower the number further, and can lead to a desired vacuum energy of order $10^{-46}\,\mathrm{GeV}^4$. In this scenario, the vacuum energy suppressed by M_P^3 can interpret the correct magnitude of the negative decceleration parameter [7].

On the other hand, there exists another plausible scale for supersymmetry breaking, the intermediate scale at $10^{10-13}\,\mathrm{GeV}$ [8]. In eq. (2), we did not consider this scale under the assumption that the real physical parameters at the observable sector are provided by the mass splitting of superpartners. In this theory, the vacuum energy and its one-loop correction are of order the intermediate scale [9], but in a theory with the vanishing cosmological constant following the philosophy of ref. [7],

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our starting point is that the vacuum energy is zero at the minimum of the potential. However, this intermediate scale must be scrutinized in specific quintessence models.

The best candidate for quintessence is a pseudo-Goldstone boson [6, 7, 10].[1] In string models, usually the string axions belong to this category [7, 10]. The key point for this argument is that in a specific direction of the parameter space the potential is almost flat to realize the extremely shallow quintessence potential. Namely, the fundamental theory must have a spontaneously broken global symmetry with a very tiny explicit breaking term of order $1/M_P^3$ [7].

In supergravity, the supersymmetry breaking is introduced by the confining force in the hidden sector [8, 12]. As usual, we introduce the hidden sector as appearing at 10^{10-13} GeV which contains a confining gauge group.

Even though the model presented in ref. [7], with a hidden sector $\mathrm{SU}(2)_H$, nicely introduces the quintessence vacuum energy, the model itself destroys the successful gauge coupling unification due to the introduction of the extra $\mathrm{SU}(2)_W$ triplet fields in the model. The $\mathrm{SU}(2)_W$ gauge group has been used in ref. [7] to protect the hidden-sector (h-sector) quarks from obtaining mass. Toward the gauge coupling unification, a gauge group needed to protect the h-sector quark from obtaining mass might be restricted to U(1)'s so that the difference of $\mathrm{SU}(3)_c$ and $\mathrm{SU}(2)_W$ β functions is not changed from the minimal supersymmetric standard model(MSSM). The β function for $\mathrm{U}(1)_Y$ can also be unchanged if the h-sector quarks carry zero electromagnetic charge. Therefore, we must decouple the hidden sector from the observable sector.

In this paper, we try to build a quintessence model purely from the theory of the hidden sector decoupled from the observable sector. The hidden sector non-abelian group can be semi-simple, but we will restrict to the simple group $\mathrm{SU}(3)_H$ as an example. The group SU(2) is smaller, but starting with the universal coupling constant at the Planck scale, SU(2) is not large enough to be strong at the intermediate scale. With $\mathrm{SU}(3)_H$, one can introduce a small number of h-quarks than in QCD and can make the coupling strong at the intermediate scale.[2] We will also introduce a U(1) factor gauge group, $\mathrm{U}(1)_H$, and the h-sector gauge group is assumed to be

$$\mathrm{SU}(3)_H \times \mathrm{U}(1)_H\,. \tag{3}$$

If the extra superfields do not carry any quantum number of $\mathrm{SU}(3)_c\times\mathrm{SU}(2)_W\times\mathrm{U}(1)_Y$, the h-sector is a real hidden sector from the observable sector. Its effect is transmitted to the observable sector only by gravity. We introduce four h-sector chiral quarks (h-quarks) and three singlets whose quantum numbers are shown in table 1.
The renormalizable terms for the singlets are

$$W = f_0\, S_0^3 + f_1 S_1 S_2 S_0 - \mu\, S_1 S_2 - \mu' S_0^2 + \mu''^2 S_0\,. \tag{4}$$

[1] For more quintessence idea, see ref. [11].

[2] If needed, one can introduce $\mathrm{SU}(4)_H$ not changing our conclusion.

	Q_1	Q_2	$\bar{Q}_3$	$\bar{Q}_4$	S_1	S_2	S_0
$\mathrm{SU}(3)_H$	$\mathbf{3}$	$\mathbf{3}$	$\mathbf{3}^*$	$\mathbf{3}^*$	$\mathbf{1}$	$\mathbf{1}$	$\mathbf{1}$
$\mathrm{U}(1)_H$	4	-4	0	0	1	-1	0

Table 1: The $\mathrm{SU}(3)_H$ and $\mathrm{U}(1)_H$ quantum numbers of superfields.

The vanishing $\mathrm{U}(1)_H$ D-term sets $|\langle S_1\rangle| = |\langle S_2\rangle|$. Vanishing of the F-term gives

$$\tilde{v}\,(-\mu + f_1 S_0) = 0\,, \qquad 3\,f_0\,S_0^2 + f_1 S_1 S_2 - 2\,\mu' S_0 + \mu''^2 = 0\,, \tag{5}$$

where we have the usual μ and μ' terms, and $|\langle S_1\rangle| = |\langle S_2\rangle| = \tilde{v}$. μ, μ' and μ'' are assumed to be at TeV scale, for which there can be many solutions [4]. These solutions assume some kind of symmetry such as $\mathrm{U}(1)_{PQ}$, $\mathrm{U}(1)_R$, or others. If the symmetry is global, one expects a Goldstone boson when it is broken. In this paper, however, we will try to introduce a global symmetry from a gauge symmetry and hence treat these mass parameters μ, μ' and μ'' just as inputs. We will comment more on the μ parameter later.

Let $\langle S_1\rangle = e^{i(\alpha+\beta)}\tilde{v}$ and $\langle S_2\rangle = e^{i(\alpha-\beta)}\tilde{v}$. Then, we obtain

$$\tilde{v} = \frac{e^{-i\alpha}\mu}{f_1}\sqrt{-\left(3\,\frac{f_0}{f_1} - 2\,\frac{\mu'}{\mu} + f_1\,\frac{\mu''^2}{\mu^2}\right)}\,. \tag{6}$$

Thus the $\mathrm{U}(1)_H$ symmetry is broken at a TeV scale and we expect a neutral gauge boson at this scale. Since this hidden sector $\mathrm{U}(1)_H$ interacts with the observable sector only through gravity, it is not expected for this extra $\mathrm{U}(1)_H$ gauge boson to be found at high energy accelerators.

The most dominant $\mathrm{SU}(3)_H \times \mathrm{U}(1)_H$ invariant couplings are

$$\frac{1}{M_P^3}\,Q_1\bar{Q}_3\,S_2^4\,, \qquad \frac{1}{M_P^3}\,Q_1\bar{Q}_4\,S_2^4\,, \qquad \frac{1}{M_P^3}\,Q_2\,\bar{Q}_3\,S_1^4\,, \qquad \frac{1}{M_P^3}\,Q_2\,\bar{Q}_4\,S_1^4\,, \tag{7}$$

which can lead to the desired vacuum energy since it is suppressed by M_P^3. The terms given in eq. (7) break the following chiral symmetry

$$\{Q_1, Q_2, \bar{Q}_3, \bar{Q}_4\} \longrightarrow e^{i\alpha}\,\{Q_1, Q_2, \bar{Q}_3, \bar{Q}_4\}\,. \tag{8}$$

The S_i fields are invariant under the above chiral transformation. The h-quarks would obtain mass of order $\tilde{v}^4/M_P^3 \sim 2.7\times 10^{-46}\,\mathrm{GeV}$ which leads to the almost zero h-sector instanton potential,

$$\sim \frac{m_Q^2\,M_{\text{h-gluino}}^3}{\Lambda_\mathrm{h}} \sim 10^{-98}\,\mathrm{GeV}^4 \tag{9}$$

for $M_{\text{h-gluino}} \sim 100\,\mathrm{GeV}$, $M_Q \sim S_i^4/M_P^3$ $(i = 1, 2)$ and $\Lambda_\mathrm{h} \sim 10^{13}\,\mathrm{GeV}$. Therefore, the instanton potential is not important, compared to the above explicit chiral symmetry breaking terms given in eq. (7). The h-squarks (the superpartner of h-quark) are assumed to have no condensation. This point is essential for interpreting a model-dependent axion $a_{\rm MD}$ as the quintessence.

First we note the following. If there is a hierarchy between two scales for two pseudoscalar bosons both of which couple to two independent explicit symmetry breaking terms, then one can easily estimate the decay constants and masses of the bosons due to the *inverted order decay constant theorem*: *The smaller decay constant corresponds to the larger explicit symmetry breaking scale, and the larger decay constant corresponds to the smaller explicit symmetry breaking scale* [13, 14, 7].

To show this, let us consider two Goldstone bosons, a_1 and a_2, with decay constants F_1 and F_2, respectively, arising from the spontaneous breaking scales F_1 and F_2 of two global symmetries. Without explicit symmetry breaking terms, there will be no potential depending on a_1 and a_2, due to the Goldstone theorem. If there exist explicit symmetry breaking terms, there result a_1- and a_2- dependent potentials which must be periodic functions of a_1 and a_2. Since the original global symmetry can be called phase symmetries, the fields must return to itself after shifts of $a_1 \to a_1 + 2\pi F_1/\alpha_1$ or $a_2 \to a_2 + 2\pi F_2/\alpha_2$ where α_1 and α_2 are determined by the quantum numbers of the fields breaking U(1) global symmetries. Therefore, if there exist two explicit symmetry breaking terms, then the potential is of the following form,

$$V \sim -V_1 \cos\left(\alpha_1 \frac{a_1}{F_1} + \alpha_2 \frac{a_2}{F_2}\right) - V_2 \cos\left(\beta_1 \frac{a_1}{F_1} + \beta_2 \frac{a_2}{F_2} + \text{constant}\right),$$

where the hierarchical explicit symmetry breaking scales are parametrized by V_1 and V_2, and both α_1 and α_2 are nonzero (i.e. two disconnected sector is not considered). To give masses to both bosons, we require $\alpha_1/\alpha_2 \neq \beta_1/\beta_2$. For a hierarchy, $V_1 \gg V_2$, one Goldstone boson determined by the dominant symmetry breaking scale is approximately,

$$\alpha_1 \frac{a_1}{F_1} + \alpha_2 \frac{a_2}{F_2} \propto \frac{a'}{F'} ,$$

where

$$F' = \frac{F_1 F_2}{\sqrt{\alpha_1^2 F_2^2 + \alpha_2^2 F_1^2}} .$$

F' is the decay constant corresponding to the mass eigenstate a'. If there is a hierarchy among the decay constants, say $F_1 \gg F_2$, then we obtain $F' \simeq F_2/\alpha_2$, i.e. the decay constant corresponding to the larger explicit symmetry breaking scale V_1 is the smaller (i.e. F_2) of F_1 and F_2, leading to the inverted order theorem. Orthogonalization of the mass matrix leads to the other mass eigenstate whose decay constant is the larger scale F_1 and explicit symmetry breaking scale is the smaller one V_2.

There are two kinds of axions in superstring models. The model-independent axion is present in any superstring models. It is basically the second rank antisymmetric tensor field $B_{\mu\nu}$ ($\mu, \nu = $ 4D indices) after compactification. The model-dependent axion couples to the fields universally. In addition, there exist model-dependent axions B_{ij} ($i, j =$ internal indices) whose properties depend on how the model is compactified. The numbers of model-independent and model-dependent axions are the zeroth and the second Betti numbers, respectively.

One very well-known example is the η-axion mixing. The chiral symmetry breaking of QCD introduces the QCD scales $\sim 1\,\mathrm{GeV}$ and $100\,\mathrm{MeV}$ as the explicit symmetry breaking scales for η and axion, respectively. The decay constants of η-meson and axion are of order $1\,\mathrm{GeV}$ and $F_a \sim 10^{9-12}\,\mathrm{GeV}$ [2], respectively. Due to the above theorem, η obtains a mass of order $1\,\mathrm{GeV}$, and axion obtains a mass of order $(0.1\,\mathrm{GeV})^2/F_a$. Another example is the axion decay constant calculated in ref. [7, 15].

In the present case, however, we have four Goldstone bosons, the model-independent axion $a_{\rm MI}$, a model dependent axion $a_{\rm MD}$, the pseudoscalar $a_{\text{h-gl}}$ arising from the h-sector gluino condensation of order $\Lambda_{\rm h}^3 \sim (10^{13}\,\mathrm{GeV})^3$, and the pseudoscalar $a_{\text{h-qu}}$ arising from the h-quark condensation below the electroweak scale.[3] Therefore, in the present case one cannot state as simply as for the case of two Goldstone bosons. Toward an explicit discussion, we will choose the decay constant scales for $a_{\rm MI}$ and $a_{\rm MD}$ as $F_a \sim 10^{12}\,\mathrm{GeV}$ and $F_q \sim M_P$, respectively [7]. The decay constant for $a_{\text{h-gl}}$ is $f_{\text{h-gl}}\Lambda_{\rm h}$, and the decay constant for $a_{\text{h-qu}}$ is about the electroweak scale, $\langle\text{h-quark}\cdot\text{h-quark}\rangle \sim (f_{\text{h-qu}}v)^3$ where $v \simeq 250\,\mathrm{GeV}$. Here, f's are dimensionless parameters.

The explicit breaking scales giving Goldstone bosons tiny masses are the h-sector gluino condensation $\Lambda_{\rm h}$, the QCD scale $\Lambda_{\rm QCD}$, $\epsilon_{\text{h-qu}} \sim v^{7/4}/M_P^{3/4}$ from eq. (7), ϵ_Λ which arises from the h-quark mass suppressed h-sector instanton potential, and $f_{\text{h-gl}}^{3/2}\Lambda_{\rm h}^{3/2}/M_P^{1/2}$. Numerically, ϵ_Λ is negligible compared to the others, and hence we will neglect ϵ_Λ. Among these, the largest explicit breaking scale is $\Lambda_{\rm h}$. The Goldstone boson affected by $\Lambda_{\rm h}$ is $a_{\text{h-gl}}$. Thus, $a_{\text{h-gl}}$ obtains a mass of order $\Lambda_{\rm h}$. Among the remaining scales, the largest explicit breaking scale is $\Lambda_{\rm QCD}$. The Goldstone boson $a_{\text{h-qu}}$ corresponding to the h-quark is separated from QCD. Thus, among the other two, $a_{\rm MI}$ and $a_{\rm MD}$, where we choose their decay constants at the scales $10^{12}\,\mathrm{GeV}$ and M_P, respectively, $a_{\rm MI}$ corresponds to $\Lambda_{\rm QCD}$ and $a_{\rm MD}$ corresponds to M_P according to the *inverted order decay constant theorem.* Thus, the very light axion for the strong CP solution can be realized by $a_{\rm MI}$ with an appropriate $F_a \sim 10^{12}\,\mathrm{GeV}$ [7, 16]. Now the shallow potential for the remaining $a_{\rm MD}$ is compared to the other explicit breaking terms and we can consider 2×2 mass matrix of the remaining Goldstone bosons, $a_{\rm MD} \simeq a_q$ and $a_{\text{h-qu}}$. These two pseudoscalars have the hidden sector interaction. The two decay constants have a hierarchy $F_q \sim M_P \gg f_{\rm h-qu}v$. Thus, again from the *inverted order decay constant theorem*, $a_{\text{h-qu}}$ with the decay constant $\sim v$ corresponds to the explicit breaking scale $\epsilon_{\text{h-qu}}$ and a_q corresponds to the explicit breaking

[3]If the fermion bilinears condense, then supersymmetry is broken. In this paper, we assume supersymmetry below the h-sector scale, and hence the h-quark condensation scale is assumed to be of order the mass splitting between superpartners.

scale $(10^{-46}\,\mathrm{GeV}^4)^{1/4} \simeq 10^{-11.5}\,\mathrm{GeV}$. Thus the mass of the Goldstone boson $a_{\text{h-qu}}$ is about $\epsilon^2_{\text{h-qu}}/f_{\text{h-qu}}v \sim 2.6 \times 10^{-22}\,\mathrm{GeV}/f_{\text{h-qu}}$, and the remaining a_q which we interpret as quintessence has the decay constant $F_q \sim M_P$ and mass $\sim 10^{-33}\,\mathrm{eV}$. Since the model-dependent axion which arises from the compactification scheme is interpreted as quintessence, it is possible to expect $F_q \sim M_P$ in some compactification models.

Now we proceed to discuss the couplings of the quintessence to matter fields. Let us assume that the quintessence is a model-dependent axion so that we use $a_{\mathrm{MD}} \simeq a_q$ with the decay constant $F_{\mathrm{MD}} \simeq F_q$. Then, we can imagine the couplings of $a_{\mathrm{MD}} \simeq a_q$, being moduli,

$$\frac{1}{M_P^3}\, e^{i a_{\mathrm{MD}}/F_q}\, Q_1\, \bar{Q}_3\, S_2^4, \ldots \tag{10}$$

Then the terms given in eq. (7) is invariant under the transformation (8) with the shift of a_{MD}

$$a_{\mathrm{MD}} \longrightarrow a_{\mathrm{MD}} - \alpha\, F_{\mathrm{MD}}\,. \tag{11}$$

If the first term in eq. (10) is the only allowed coupling, then the supersymmetric F-terms do not generate a potential for a_{MD}. However, there can exist other terms denoted by dots. In this case, the supersymmetric terms can generate a_{MD} dependent potential; but these are suppressed by M_P^6, viz. $V \sim \sum_i |\partial W/\partial\phi_i|^2$. With a broken supersymmetry, the soft terms, i.e. $m_{3/2} \cdot$ (eq. (10)), can give a potential depending on a_q. This is a situation we imagine to be realized for a natual quintessence. The vacuum energy generated by $\langle Q_1 \bar{Q}_3\, S_2^4\rangle$ breaks the above shift symmetry, and the Goldstone boson corresponding to the shift symmetry obtains a mass of order $(\langle V\rangle/F_q^2)^{1/2} \sim 10^{-33}\,\mathrm{eV}$. For this argument to be valid, the superpotential terms are restricted by some kind of symmetry (using global and discrete symmetries) so that superpotential terms with $D \leq 5$ are forbidden.

But in string models, the global symmetries are badly broken by the world-sheet instanton effect [17, 18]. The coupling of the model-dependent axion is given by

$$W \sim \Phi_1\Phi_2\Phi_3\left(1 + \epsilon\, e^{-R/M_P}\right), \tag{12}$$

where ϵ denotes the importance of the world-sheet contribution, $R = r - i\, a_{\mathrm{MD}}$ is the chiral field containing a_{MD}, and $\Phi_1\Phi_2\Phi_3$ represents a generic superpotential term. As discussed in the previous paragraph, coupling of the model-dependent axion to superpotential terms with $D \leq 5$ can be too large for a successful identification of a_{MD} as a_q. But there is a caveat in this argument. The dangerous a_{MD} coupling in eq. (12) for it to be quintessence arises if the vacuum expectation value $\langle\Phi^3\rangle$ is much largerthan $(0.003\,\mathrm{eV})^4$.

If the terms in the superpotential do not generate non-vanishing VEV's, then the model-dependent axion coupling is negligible. From the superpotential of the type (12), one obtains soft terms and fermionic terms of the following form:

$$V_{\rm soft} \sim m_{3/2}\,\Phi_1\Phi_2\,\Phi_3\,\left(3+3\epsilon e^{-R/M_P}+\epsilon e^{-2R/M_P}\right)+\text{h.c.}\,, \tag{13}$$

$$-\mathcal{L}_Y \sim -\psi_{\Phi_1}\psi_{\Phi_2}\Phi_3\left(1+\epsilon\, e^{-R/M_P}\right)+\text{h.c.}+(\text{permutations})+(\text{axino terms})\,, \tag{14}$$

where ψ_{Φ_i} field is the fermionic superpartner of Φ_i.

Interactions are obtained from the superpotential (12) where we suppressed couplings for simplicity. The potential

$$V=\sum_{\phi_i=\Phi_1,\Phi_2,\Phi_3,R}\left|\frac{\partial W}{\partial\phi_i}\right|^2 \tag{15}$$

includes the $a_{\rm MD}$-dependent terms of the following form,

$$\left|\Phi_1\Phi_2\left(1+\epsilon\, e^{-R/M_P}\right)\right|^2+(\text{permutations})\,, \tag{16}$$

from which we extract dominant supersymmetric $a_{\rm MD}$-dependent couplings,

$$2\,\epsilon\, e^{-r/M_P}\,|\Phi_1\Phi_2|^2\cos\frac{a_{\rm MD}}{M_P}+(\text{permutations})\,, \tag{17}$$

where we shifted $a_{\rm MD}$ so that ϵ becomes real.

From the above example, one can see that the most dangerous terms can come from superpotentials constructed from SU(3) × SU(2) × U(1) singlet fields. If all the singlets involved in the couplings, viz. eqs. (15) and (16), obtain superheavy vacuum expectation values, our argument for the quintessence cannot be realized. Therefore, we assume that superpotentials constructed with singlet fields do not give non-vanishing vacuum energy, namely at least one scalar component of Φ_i in eq. (13) has the vanishing VEV. For eq. (16) not to contribute, at most one out of Φ_1,Φ_2 and Φ_3 can have a nonvanishing VEV. Note that the fermionic components of the singlets do not condense and eq. (14) is not particularly important for the singlets.

But eq. (14) can be important for the fields of the minimal supersymmetric standard model(MSSM). In the MSSM let us first consider $V_{\rm soft}$. The A-terms of the MSSM Yukawa couplings in the charge and color conserving vacuum do not generate a $a_{\rm MD}$ dependent potential. Among the MSSM couplings, therefore, the relevant term is the B-term from the μ-term [4] since $\langle H_1H_2\rangle\neq 0$. If the μ term is corrected as above by the world-sheet instanton effect, then the model-dependent axion cannot be identified as the quintessence. However, the argument given in ref. [18] does not apply to the μ-term.[4]

[4]In ref. [18], $\Phi_1\Phi_2\Phi_3e^{-R}$ coupling was obtained by considering the expectation value of the dilaton vertex operator V_B between the fermionic states, e.g. $\langle\frac12,\frac12,\frac12,-\frac12\,|V_B|-\frac12,-\frac12,-\frac12,\frac12\rangle$. For the μ term, however, one needs to calculate $\langle\frac12,\frac12,\frac12,-\frac12\,|1|-\frac12,-\frac12,-\frac12,\frac12\rangle$ which seems to be vanishing.

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Next, the most dangerous term is the Yukawa couplings $\mathcal{L}_Y$. The known quark condensation and VEV of the Higgs doublet fields would generate an electroweak scale vacuum energy which is enormous compared to the quintessence potential if it is coupled to $a_{\rm MD}$. But if the model-dependent axion potential to the MSSM fields is *universal*, then the Yukawa coupling is not dangerous. In this case, we can write the coupling as

$$-\left(1+\epsilon\, e^{-R}\right)\cdot\left(\sum_{ij} f^d_{ij}\,\bar{d}^i q^j H_1+\sum_{ij} f^u_{ij}\,\bar{u}^i q^j H_2\right)+\text{h.c.}\,, \tag{18}$$

where i, j are the family indices. At the minimum of the potential the vacuum energy is zero, i.e.

$$-\left\langle\sum_{ij} f^d_{ij}\,\bar{d}^i q^j H_1+\sum_{ij} f^u_{ij}\,\bar{u}^i q^j H_2\right\rangle=0\,. \tag{19}$$

Therefore, at the minimum the tree level $a_{\rm MD}$ potential is vanishing even if we consider the quark condensation.

At the electroweak scale another relevant term in the superpotential is the μ term, $\mu\, H_1 H_2$. If we treat μ as an input parameter, then we can neglect its coupling to $a_{\rm MD}$ through eq. (12), as commented above. But in some solutions of the μ problem, it is generated through non-renormalizable terms in which case we should consider the coupling of the type eq. (12). Then it leads to a too large potential for the model-dependent axion to be the quintessence. For example, one may consider only cubic terms in the superpotential W_0 inspired by string models [20]. Therefore, $\mu\, H_1 H_2$ is absent in W_0. W_0 is assumed to contain the singlet terms with the intermediate scale VEV's. Supergravity can contain non-renormalizable terms $W_0\, H_1 H_2/M_P^2$ which will give the desired magnitude of μ for $\langle W_0\rangle\sim$ (intermediate scale)^3. For this non-renormalizable term, we expect that there exists a world-sheet instanton correction which is not expected to take the universal form eq. (12). In this case, $a_{\rm MD}$ cannot work as the quintessence. In the same vein, any solution of μ by the superpotential is thought to have the same problem. However, if the world-sheet instanton correction to the non-renormalizable potential also takes the universal form eq. (12), then the requirement of vanishing vacuum energy gives the $a_{\rm MD}$ coupling negligible.

In conclusion, we investigated a model to interpret a model-dependent axion $a_{\rm MD}$ as quintessence. For this idea to work, there must exist a symmetry so that superpotential terms generating a vacuum energy of order $v_{\rm ew}^6/M_P^2$ must be forbidden. We realize this idea introducing a $\mathrm{SU}(3)\times\mathrm{U}(1)_H$ gauge symmetry in the hidden sector.

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References

[1] S. Weinberg, *Implications of dynamical symmetry breaking: an addendum*, *Phys. Rev.* **D 19** (1979) 1277;
L. Susskind, *Dynamics of spontaneous symmetry breaking in the Weinberg-Salam theory*, *Phys. Rev.* **D 20** (1979) 2619;
E. Witten, *Dynamical breaking of supersymmetry*, *Nucl. Phys.* **B 188** (1981) 513.

[2] J.E. Kim, *Light pseudoscalars, particle physics and cosmology*, *Phys. Rep.* **150** (1987) 1;
H.-Y. Cheng, *The strong CP problem revisited*, *Phys. Rep.* **158** (1988) 1;
R.D. Peccei, *The strong CP problem*, in *CP violation*, C. Jarlskog ed., World Scientific, 1989, p.503.

[3] J.E. Kim, *Constraints on very light axions from cavity experiments*, *Phys. Rev.* **D 58** (1998) 055006 [hep-ph/9802220];
C. Hagmann et al., *Results from a high-sensitivity search for cosmic axions*, *Phys. Rev. Lett.* **80** (1998) 2043 [astro-ph/9801286].

[4] J.E. Kim and H.P. Nilles, *The μ problem and the strong CP problem*, *Phys. Lett.* **B 138** (1984) 150;
G.F. Giudice and A. Masiero, *A natural solution to the μ problem in supergravity theories*, *Phys. Lett.* **B 206** (1988) 480;
J.A. Casas and C. Muñoz, *A natural solution to the μ problem*, *Phys. Lett.* **B 306** (1993) 288 [hep-ph/9302227];
J.E. Kim and H.P. Nilles, *Symmetry principles toward solutions of the μ problem*, *Mod. Phys. Lett.* **A 9** (1994) 3575 [hep-ph/9406296].

[5] S. Perlmutter et al., *Measurements of the cosmological parameters Ω and Λ from the first 7 supernovae at $z \geq 0.35$*, *Astrophys. J.* **483** (1997) 565 [astro-ph/9608192];
The Supernova Cosmology Project, *Cosmology from type-IA supernovae*, *Bull. Am. Astron. Soc.* **5** (1998) 1351 [astro-ph/9812473];
A. Riess, W.H. Press and R.P. Kirshner, *Using SN-IA light curve shapes to measure the Hubble constant*, submitted to *Astrophys. J.* (1998);
S. Perlmutter and A. Riess, *Cosmological parameters from supernovae: two group's results agree*, in *COSMO-98*, D. Caldwell ed., AIP Conference Proc. 478, Woodbury, N.Y., 1999, p. 129.

[6] J. Frieman, *Dynamical vacuum*, talk presented at COSMO-98, Asilomar, California, Nov. 15–20, 1998 (unpublished).

[7] J.E. Kim, *Axion and almost massless quark as ingredients of quintessence*, *J. High Energy Phys.* **05** (1999) 022 [hep-ph/9811509].

[8] H.P. Nilles, *Supersymmetry, supergravity and particle physics*, *Phys. Rep.* **110** (1984) 1.

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[9] K. Choi, J.E. Kim and H.P. Nilles, *Cosmological constant and soft terms in supergravity*, *Phys. Rev. Lett.* **73** (1994) 1758 [hep-ph/9404311].

[10] K. Choi, *String or M-theory axion as a quintessence*, hep-ph/9902292.

[11] P. Binetruy, *Models of dynamical supersymmetry breaking and quintessence*, *Phys. Rev.* **D 60** (1999) 063502 [hep-ph/9810553];
C. Kolda and D.H. Lyth, *Quintessential difficulties*, *Phys. Lett.* **B 458** (1999) 197 [hep-ph/9811375];
T. Chiba, *Quintessence, the gravitational constant and gravity*, *Phys. Rev.* **D 60** (1999) 083508 [gr-qc/9903094];
P. Brax and J. Martin, *Quintessence and supergravity*, *Phys. Lett.* **B 468** (1999) 40 [astro-ph/9905040];
A. Masiero, M. Pietroni and F. Rosati, *Susy QCD and quintessence*, *Phys. Rev.* **D 61** (2000) 023504 [hep-ph/9905346];
M.C. Bento and O. Bertolami, *Compactification, vacuum energy and quintessence*, *Gen. Rel. Grav.* **31** (1999) 1461 [gr-qc/9905075];
F. Perrotta, C. Baccigalupi and S. Matarrese, *Extended quintessence*, *Phys. Rev.* **D 61** (2000) 023507 [astro-ph/9906066].

[12] J.P. Derendinger, L.E. Ibanez and H.P. Nilles, *On the low-energy $D = 4$, $N = 1$ supergravity theory extracted from the $D = 10$, $N = 1$ superstring*, *Phys. Lett.* **B 155** (1985) 65;
M. Dine, R. Rohm, N. Seiberg and E. Witten, *Gluino condensation in superstring models*, *Phys. Lett.* **B 156** (1985) 55.

[13] K. Choi and J.E. Kim, *Compactification and axions in $E_8 \times E_8'$ superstring models*, *Phys. Lett.* **B 165** (1985) 71.

[14] E.J. Chun, J.E. Kim and H.P. Nilles, *A natural solution of the mu problem with a composite axion in the hidden sector*, *Nucl. Phys.* **B 370** (1992) 105.

[15] H. Georgi, J.E. Kim and H.P. Nilles, *Hidden sector gaugino condensation and the model independent axion*, *Phys. Lett.* **B 437** (1998) 325 [hep-ph/9805510];
J.E. Kim, *Superstring axion, gaugino condensation and discrete symmetries*, in *PASCOS-98*, P. Nath ed., World Scientific, Singapore 1999, p. 361, talk presented at PASCOS-98, Boston, March 22-29, 1998 [hep-ph/9807322].

[16] J.E. Kim, *Weak interaction singlet and strong CP invariance*, *Phys. Rev. Lett.* **43** (1979) 103;
M.A. Shifman, V.I. Vainstein and V.I. Zakharov, *Instantons in nonperturbative QCD vacuum*, *Nucl. Phys.* **B 165** (1980) 45;
M. Dine, W. Fischler and M. Srednicki, *A simple solution to the strong CP problem with a harmless axion*, *Phys. Lett.* **B 104** (1981) 199;
A. Zhitniskii, *On possible suppression of the axion hadron interactions*, (in Russian), *Sov. J. Nucl. Phys.* **31** (1980) 260.

JHEP06(2000)016

[17] X.G. Wen and E. Witten, *World sheet instantons and the Peccei-Quinn symmetry*, *Phys. Lett.* **B 166** (1986) 397.

[18] M. Dine, N. Seiberg, X.G. Wen and E. Witten, *Nonperturbative effects on the string world sheet*, *Nucl. Phys.* **B 278** (1986) 769.

[19] J.E. Kim, *A natural solution of μ from the hidden sector*, *Phys. Lett.* **B 452** (1999) 255 [hep-ph/9901204].

[20] J.A. Casas and C. Muñoz, *A natural solution to the μ problem*, *Phys. Lett.* **B 306** (1993) 288 [hep-ph/9302227].