QCD Axion and Quintessential Axion

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Abstract. The axion solution of the strong CP problem is reviewed together with the other strong CP solutions. We also point out the quintessential axion (quintaxion) whose potential can be extremely flat due to the tiny ratio of the hidden sector quark mass and the intermediate hidden sector scale. The quintaxion candidates are supposed to be the string theory axions, the model independent or the model dependent axions.

1 The Strong CP Problem

At this conference, the organizer asked to review on the axion [1] as the CDM candidate [2]. Therefore, I will review the QCD axion and the related problem, the strong CP problem. With supersymmetrization, an O(GeV) axino can be a CDM candidate also [3] in addition to the axion CDM, but we will not discuss this interesting class of CDM candidate here. It was reviewed in another DM conference [4]. Out of several CDM candidates, the axion is most attractive since it arises from the need to solve the strong CP problem. One crucial difference of the axion CDM from most other CDM candidates is that the axion is a classical field oscillation [2] while many other CDM candidates rely on their heavy particle nature. The classical axion potential is extremely flat as shown in Fig. 1. The axion value \( \langle a \rangle \) stays at some point (the bullet) for a long time before it starts to oscillates, which occurs at \( T \simeq 1 \text{ GeV} \) when the Hubble time \( 1/H \) becomes larger than the oscillation period \( m_a \).

![Fig. 1. The very flat axion potential. The minimum of the potential is shown by a triangle. The initial axion vacuum is shown as a bullet.](image)

This axion is the invention from the need to solve the strong CP problem. The instanton solution in nonabelian gauge theories needs an effective term parametrized by \( \xi \)

\[
\frac{\bar{\theta}}{32\pi^2} F \tilde{F} \tag{1}
\]

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where $\tilde{F} = \epsilon_{\mu
u\rho\sigma} F^{\mu\nu} F^{\rho\sigma}$ is the Pontryagin term of the gluon field strength $F^a_{\mu\nu}$. The $\bar{\theta}$ parameter is the final low energy value, taking into account the weak CP violation. It can be split into two pieces $\bar{\theta} = \theta_{QCD} + \theta_{weak}$ where $\theta_{QCD}$ is the parameter determined at high energy scale and $\theta_{weak} = \text{Arg} \text{Det} M_q$ is the one generated at the electroweak scale. Since $\bar{\theta}$ contributes to the neutron electric dipole moment (NEDM), the experimental upper bound on NEDM constrains $|\bar{\theta}| < 10^{-9}$, which is a fine-tuning problem: the so-called strong CP problem.

One may argue that there is no strong CP problem from the beginning. One such argument is a 5D extension where there is no 5D instanton solution \cite{6}. However, I think that in the 4D effective theory language where 5D is hidden one should consider the $\bar{\theta} F \tilde{F}$ term at low energy, and the strong CP problem reappears.

So, in this talk I will review the solutions of the strong CP problem first on,

- the calculable $\theta_{weak}$,
- the massless up quark solution, and
- the very light axion,

and then go on to discuss an axion-like particle,

- the quintessential axion \cite{7}.

The calculable $\theta_{weak}$ and the massless up quark cases are discussed in the subsequent subsections, and the very light axion and quintessential axion in separate sections.

### 1.1 Calculable $\theta_{weak}$

In this kind of theories $\theta_{QCD}$ is supposed to be zero, which can be implemented by assuming the CP invariance of the Lagrangian. Thus, the weak CP violation can be introduced only spontaneously \cite{8}. After the weak CP violation is introduced spontaneously, $\theta_{weak}$ can be calculated. In the beginning, this kind of models were proposed by changing the weak interactions \cite{9}. However, the Kobayashi-Maskawa (KM) type weak CP violation seems to be working, and the above spontaneous weak CP violations are not working. In this regards, the Nelson-Barr approach \cite{10} of the spontaneous CP violation at a very high energy scale with a calculable $\theta_{weak}$ has attracted a great deal of attention since at low energy it leads to the KM type weak CP violation. However, it needs superheavy particles with a vectorlike representation at high energy scale.

### 1.2 Massless up quark

Suppose that we chiral-transform a quark field $q$,

$$ q \rightarrow e^{i\gamma_5 \alpha} q $$

It is equivalent to changing $\bar{\theta} \rightarrow \bar{\theta} - 2\alpha$. Thus, if a theory admits such a transformation, i.e. if there exists a massless quark, then the strong CP problem is not
present. The up quark is most promising for the massless quark possibility. It was known from the very early days of the $\theta$ vacuum that a massless up quark solves the strong CP problem. The current debate on the massless up quark solution is a phenomenological one. The famous up/down ratio $Z = m_u/m_d = \frac{1}{2}$ seems to rule out the $m_u = 0$ possibility.

But the story below 100 GeV is not so simple. The point is that there exists the 't Hooft determinental interaction shown in Fig. 2.

![Fig. 2. The 't Hooft determinental interaction is shown as arrows. The s quark line can be closed to give an effective instanton generated $\frac{m_s}{A}u_R u_L d_R d_L$ interaction.](image)

The hard quark mass considered in the chiral perturbation theory is hard at the 300 MeV scale. If the instanton generated mass is defined much above 300 MeV, then it will act like a hard mass in the chiral perturbation theory. Indeed, by closing the strange and down quarks in Fig. 2, one can obtain an up quark mass of order

$$m_u = \frac{m_s m_d}{A}$$  \hspace{1cm} (3)

where $A$ is the instanton related QCD parameter of order GeV. Suppose that the electroweak symmetry breaking gives a massless up quark. Then, there is no strong CP problem. But the chiral perturbation calculation can be performed with the above instanton generated nonzero up quark mass. So, the massless up quark problem is whether the instanton generated up quark mass is consistent with the phenomenology of the chiral perturbation. Here, there are two groups sharing the opposite view:

1 By connecting the $u_R, s_R$ quark lines of Fig. 2 with $u_L, s_L$ lines in the $G_F$ order weak interaction, the instanton generated $\Delta I = \frac{1}{2}$ was obtained long time ago [17].
Kaplan-Manohar, Choi: Yes, the up quark can be massless. Leutwyler: No, the up quark cannot be massless. Kaplan and Manohar \[11\] considered the second order chiral perturbation theory and found that \(Z = \frac{m_u}{m_d} \sim 0.2\), and concluded that \(m_u = 0\) is not ruled out. In fact, they considered the \(L_7\) term in the chiral perturbation theory, \(L_7 = \chi^T U - \chi U^T\) where \(\chi\) is a 3 \(\times\) 3 mass matrix and \(U\) is a 3 \(\times\) 3 unitary matrix in terms of the octet pseudoscalar mesons, and obtained for the case of \(m_u = 0\)

\[
L_7 \simeq (1 \sim 2) \times 10^{-4} \quad \text{[or \(2L_8 - L_5\)]}_{m_u=0} \simeq -(1.2 \sim 2.6) \times 10^{-3}
\]  

(4)

where \(L_8\) and \(L_5\) are other terms in the chiral perturbation. Leutwyler \[13\] attempted to compute \(L_7\) using the QCD sum rule with the \(\eta'\) (the SU(3) singlet meson) dominance, in analogy with the vector dominance model, and obtained \(L_7 \simeq -(2 \sim 4) \times 10^{-4}\). If Leutwyler’s calculation is correct, the \(m_u = 0\) possibility is ruled out. However, there are arguments con Leutwyler. For example, the \(\eta'\) dominance can be questioned. To answer this problem, Choi \[14\] followed the modern interpretation for the \(\eta'\) mass, i.e. the instanton contribution to the \(\eta'\) mass, in which case he showed that he could change the sign of \(L_7\), to \(L_7 \simeq (3 \sim 8) \times 10^{-4}\). Choi’s result is obtained by semiclassical approximation supplemented by gluon condensate which may be the only unreliable point. Anyway, it was noted that instanton effects suppress \(\eta'\) contribution to \(L_7\). So, it is fair to say that the problem is not completely settled in the chiral perturbation theory.

There has been a new development in the lattice calculation of the chiral perturbation parameter. For example, a lattice calculation \[18\] gives \[19\]

\[
2L_8 - L_5 \simeq (9 \pm 4_{\text{stat}} \pm 20_{\text{syst}}) \times 10^{-5}
\]  

(5)

with \(m_u/m_d \simeq 0.484 \pm 0.027\). If it is true, the \(m_u = 0\) possibility is closed. However, there can be questions on the lattice calculation such as the small number of lattice sites and no inclusion of the instanton effects. Anyway, the possibility of \(m_u = 0\) is raised from the instanton contribution to the up quark mass. Therefore, it is fair to say that the \(m_u = 0\) possibility is not closed yet from lattice calculation.

2 The QCD Axion

Axion was noted \[20\] from the spontaneously broken Peccei-Quinn symmetry \[21\]. More generally, we can define an axion \(a\) as the pseudoscalar WITHOUT ANY POTENTIAL EXCEPT that arising from the following nonrenormalizable coupling,

\[
\frac{a}{F_a} \left\{ F F \right\} \equiv \frac{1}{32\pi^2} \frac{a}{F_a^2} \epsilon_{\mu \nu \rho \sigma} F^{\mu \nu} F^{\rho \sigma}
\]  

(6)

where \(F_a\) is called “axion decay constant”. As shown in Table 1, this nonrenormalizable interaction can arise in several different ways \[22\] \[25\] \[21\] \[27\] \[28\].

The axion potential arising after integrating out the gluon field has the following two properties,
Table 1. Nature of the axion

| Source of the non-renormalizable interaction | \( F_a \sim \) |
|---------------------------------------------|-----------------|
| Higher dimensional fundamental interaction  | Planck scale \( M_P \) |
| In composite models                         | Compositeness scale \( M_c \) |
| Spontaneously broken PQ symmetry            | PQ symmetry breaking scale \( \tilde{v} \) |

**Fig. 3.** The minimum of the free energy of \( \bar{\theta} \) parameter is at \( \bar{\theta} = 0 \).

- It is periodic with the periodicity \( 2\pi F_a \),
- Its minimum is at \( \langle a \rangle = 2n\pi F_a \) (\( n \) = integer) \( [21, 29] \).

Taking into account the above points, one can write a cosine potential of \( a: -A_{QCD}^0 \cos(a/F_a) \). The height of the potential is determined by the strong interaction scale \( \Lambda_{QCD} \). For the vanishing cosmological constant, we can write the potential as \( A_{QCD}^0(1 - \cos \frac{a}{F_a}) \), where \( A_{QCD}^0 \) is an educated guess. Actually, \( A_{QCD}^0 \) must be qualified by considering the axion mixing with pions. In the current algebra or in the chiral perturbation calculation, one can fix this coefficient to obtain \( [30] \).

\[
V[a] = \frac{Z}{(1 + Z)^2} f^2 \frac{m^2}{\pi} \left( 1 - \cos \frac{a}{F_a} \right).
\]  

(7)

As we know that a massless quark makes the axion potential flat, we can guess that the coefficient is \( m_u m_d A_{QCD}^2 \) (or \( m_u A_{QCD}^3 \) at low energy) which turns out

\[ V[a] = f^2 \frac{m^2}{\pi} \left( 1 - \sqrt{m^2 + m^2} + 2m_u m_d \cos(a/F_a)/(m_u + m_d) \right). \]  

The expansion up to the \( \cos(a/F_a) \) term gives the result \( [7] \).
to be \((\text{number}) \times f_\pi^2 m_\pi^2\). We have to remember that the overall coefficient has the small quark masses, which will be very useful in the discussion of quintaxion.

The shape of the potential is shown in Fig. 3.

It shows that at the minimum of the axion potential \(\bar{\theta} = 0\), solving the strong CP problem. Basically, the axion solution of the strong CP problem is a cosmological solution. Namely, in the evolving universe, the classical axion field seeks the minimum where

\[
\bar{\theta} = 0,
\]

whatever happened before. But note that the weak interaction violates the CP invariance and introduces a small CP odd term in the axion potential, which shifts the \(\bar{\theta}\) minimum slightly to \(\bar{\theta} \sim 10^{-17}\).\(^{31}\)

From (7), one can easily compute the very light axion mass

\[
m_a = \sqrt{\frac{Z}{1+Z}} \frac{f_\pi m_\pi a}{F_a} \simeq 0.6 \, [\text{eV}] \times 10^7 \, \text{GeV} / F_a
\]

From the early days of axion physics, the PQWW axion seemed to be in conflict with laboratory experiments\(^{32}\), which was the reason that a number of calculable \(\theta_{\text{weak}}\) models were tried in 1978\(^{9}\). From the cosmological point of view, the PQWW axion is not of much significance since its lifetime is around \(10^{-8}\) s order. This has led to the invention of the invisible axions\(^{27,28}\), but it is proper to call them as very light axions since in the Sikivie type cavity experiments\(^{23}\) one may detect the Galactic axions.

The very light axions are restricted by several laboratory experiments. These include

1. meson decays: \(J/\psi \rightarrow \gamma a, \nu \rightarrow \gamma a, K^+ \rightarrow \gamma a, \pi^+ \rightarrow e^+\nu_e a, \) etc.
2. beam dump experiments: \(p(e) + N \rightarrow a + X, a \rightarrow \gamma\gamma, \) or \(e^+e^-\),
3. nuclear deexcitation: \(N^* \rightarrow Na, a \rightarrow \gamma\gamma, \) or \(e^+e^-\).

The above laboratory experiments restrict the decay constant of the very light axion to

\[
F_a > 10^4 \, \text{GeV}.
\]

However, the laboratory bound is not as strong as the supernova bound which is discussed in the following section.

3 The Axion Window from Outer Space

3.1 Stars

A stringent lower bound on the axion decay constant comes from the study of stellar objects such as Sun, red giant stars, and supernovae\(^{1}\).

In the hot plasma in stars, there exist high energy particles: \(\gamma, e, p\), and even gluons and quarks if the interior temperature is high enough. The hot plasma of the interior of stars, even axions can be produced. But the axion interaction rate is suppressed by \(F_a\). Bigger stars seemed to give stronger constraint. For example,
the red giant gave a stronger constraint than Sun, and the supernovae gave even more stronger constraint than red giant stars. Even before the discovery of SN1987A, there was a prophetic paper on supernovae by Iwamoto, and furthermore most stars are surveyed by Pantziris and Kang. In this study, the red giant constraint was the strongest, giving a bound \( F_a > 0.8 \times 10^8 \) GeV. Unfortunately, Pantziris and Kang did not give a useful bound from supernovae, presumably from the ill-treated pion-nucleon couplings.

The lucky observation of SN1987A gave the strongest lower limit on \( F_a \), \( F_a > 0.6 \times 10^9 \) GeV. In the actual calculation, a correct treatment of the interactions is needed. For example, for the pion-nucleon coupling one can use pseudo-vector derivative couplings and pseudo-scalar couplings. The reason is that the most conspicuous process of the axion emission in supernovae is

\[
N + N \rightarrow N + N + a
\]

with a pion exchange diagram. Without a correct treatment of this \( NN\pi \) coupling, the calculations cannot be trusted. It was only clarified after the initial calculations, but the bound still remains at order \( 10^9 \) GeV.

### 3.2 Cosmology

The axion has a few interesting properties relevant for cosmology. Since the axion is created by the spontaneous breaking of a global \( U(1) \) symmetry, the cosmic strings can be created in the early universe. The oscillation of this axionic string can be damped by creating axions. This was used to restrict the bound on \( F_a \), to less than \( 10^{11} \) GeV. Another interesting object created by axion background is the axionic domain walls, which can be problematic in the standard Big Bang cosmology. However, these problems are avoided in the most popular cosmological scenario, i.e. inflation with supergravity Lagrangian. Supersymmetry seems to be the most popular extension beyond the standard model, and the inflationary idea seems to be needed for the flatness and homogeneity problems. In this case, the gravitino cosmology is inescapable and the reheating temperature after inflation is better to be less than \( 10^9 \) GeV. Thus, the problematic axions from axionic strings and axionic domain walls are all inflated away.

After the inflation, the value of the classical axion field in our homogeneous patch is of order \( F_a \). For \( \langle a \rangle \ll F_a \), the probability is so small. Therefore, in general this initial value of the classical axion field is supposed to be of order \( F_a \), \( \langle a \rangle_{RH} \equiv F_1 \sim O(F_a) \). This initial value of \( \langle a \rangle \) stays there for a long time since the axion potential is extremely flat as shown in Fig. 1. The axion vacuum starts to oscillate when the Hubble expansion rate \((1/H)\) became small enough to be comparable as the classical axion oscillation rate \((1/m_a)\), which happens when the temperature of the universe falls to around \( T_1 \approx 1 \) GeV. Below \( T_1 \) the oscillating classical axion field carries the energy density of order \( m_a^2 \langle a \rangle^2 \). As the universe expands, the amplitude of the axion field shrinks, which is the result of the conserving axion number density \( m_a \langle a \rangle^2 \) in the comoving volume.
Fig. 4. The status of the axion search experiments of Rochester-Brookhaven, Florida, LLNL, Kyoto, and MIT-Fermilab experiments.
Axions

Taking into account of this evolving axion universe, one obtains

$$\rho_a(T_\gamma) = m_a(T_\gamma) n_a(T_\gamma) = \frac{2.5}{M_p} \left( \frac{\langle a \rangle(T_1)}{F_a} \right)^3 . \quad (12)$$

In Eq. (12), the factors except $\frac{F_a}{M_p}$ gives roughly $10^7 \rho_c$ where $\rho_c$ is the current critical energy density. Therefore, not to overclose the universe by the oscillating axion field, $F_a$ should be less than $10^{12}$ GeV.

Summarizing the astro- and cosmological- bounds, we arrive at an axion window still open to be observed,

$$10^9 \text{ GeV} < F_a < 10^{12} \text{ GeV}. \quad (13)$$

Note, however, that if the axions constitute only 25% of the critical energy density then $F_a$ should be $\frac{1}{4}$ of the above estimate. However, we will take $10^{12}$ GeV as a generic number since there can exist other uncertainties such as the initial value of $\langle a \rangle$, the domain wall number, etc.

3.3 Galactic axion search

If the axion is the CDM component of the universe, it is pointed out that they can be detected [33]. Even though the interaction is inversely proportional to $F_a$, the enormous number density cancels this effect, and the galactic axions can be detected. Indeed, there have been several experiments since late 1980’s [42,43]. The cavity [42] and Rydberg atom [43] searches utilize the axion-$F\tilde{F}(\text{photon})$ coupling. Even though the $\theta F\tilde{F}(\text{photon})$ is absent since $\theta$ is a parameter, the axion-$F\tilde{F}$ is a physically observable operator [33]. Sikivie considered this axion-$F\tilde{F}$ coupling in the presence of strong gradient of magnetic field of order 10 Tesla.

In Fig. 4, we summarize the current status of the axion search experiments. There are two points to be clarified here. One is that there are many very light axion models, in KSVZ type and in DFSZ type. The definition of the notation is given in Kim [44]. For example, an arbitrary choice of $e_3 = 1$ cannot fulfill the unification of couplings, and hence should not be considered toward unification of elementary particle forces [45]. Another point is that these estimates are based on closing the universe by axions. If it is required that only 25% of the critical energy is closed by the axions, the constraint on the axion density is strengthened by a factor of 4.

It is noted that still there is a lot of room to allow the very light axion as CDM.

4 The Quintaxion (Quintessential Axion)

Tyna 1a supernova data and WMAP data hints that 73% of the energy of the universe is homogeneous dark energy [40], and hence does not contribute the
galaxy formation unlike the CDM candidates. In this section, we try to introduce an axionlike particle for the dark energy, which is called a quintessential axion or quintaxion. As this name suggest quintaxion is a kind of pseudo Goldstone boson whose potential arises from instanton effects. This kind of extremely light \((\sim 10^{-33} \text{ eV})\) pseudo-Goldstone boson was suggested sometime ago \(\text{[47]}\). The attempt to incorporate it in superstring or M-theory was also tried \(\text{[48]}\). There have been another kind of scalars whose present vacuum energy is required to be \(\sim (0.003 \text{ eV})^4\), which are generally called quintessence \(\text{[49]}\).

The condition for quintaxion is \(\omega = p/\rho < -\frac{1}{3}\),\(^3\) For a pseudo-Goldstone boson \(a_q = \theta f\) with the potential \(U(\theta)\) and the decay constant \(f\), \(\omega\) is given as

\[
\omega = \frac{-6f^2 + M_p^2 |U'|^2}{6f^2 + M_p^2 |U'|^2}.
\]

(14)

With \(f\) near the Planck scale a quintaxion is naturally realized \(\text{[47]}\). For superstring axions, the decay constants are expected to be of order Planck scale but the model-independent axion decay constant is of order 10\(^{16}\) GeV \(\text{[50]}\). Except this special case, we generally anticipate that string axions such as the model-dependent ones would have a Planck scale axion decay constants. So, even though we present formulae for the model-independent axion for a concrete presentation, this caveat is implied. For this quintessential axion to work, we require that the potential is extremely flat with the height of order \((0.003 \text{ eV})^4\) with a Planck scale order decay constant.

From our discussion on the massless quark in Sec. 1, it is very tempting to use some kind of almost massless quarks for the extremely flat quintessential potential. Since a hidden sector is needed for supersymmetry breaking, we will attempt the hidden-sector quarks(h-quarks) as fulfilling this purpose. Before presenting a quintaxion along this line, let me comment a few words on the superstring axion.

Among several possibilities of the QCD axion given in Table 1, the superstring axion seems to be the most attractive realization. The axion is definitely the most attractive solution of the strong CP problem, because it arises by automatically solving the strong CP problem. It relies only on the \(aF\bar{F}\) coupling. On the other hand, the other solutions should massage the theory much, by postulating specific forms of couplings. For example, \(m_u = 0\) possibility is not that simple theoretically. It requires that Det. \(M_{up} = 0\), which is a highly non-trivial requirement for the nine complex entries of \(M_{up}\). Similarly, superstring axions is the most attractive since all string compactifications render a plenty of axions. For the other axions, one massage the theory more compared to the string axions. In this regards, if the model-independent axion were acceptable as the QCD axion, then it might have given a very reliable axion candidate.

However, the model-independent axion failed miserably because of a too large axion decay constant \(\text{[50]}\). So, as the first principle the string axion as the QCD axion failed. One must remove this model-independent axion at high energy.

\(^3\) But the recent WMAP data \(\text{[46]}\) suggests a stronger \(\omega < -0.8\).
Indeed, it is possible to do so, if there exists the anomalous $U(1)$ gauge symmetry so that the $U(1)$ gauge boson eats up the model-independent axion as its longitudinal degree [51]. But, then there must survive a global symmetry which must be broken at the intermediate scale [52]. In another context, some discrete symmetries were invoked to obtain an approximate global symmetry, approximate in the sense that the global symmetry breaking operators are allowed at dimension 13 and higher [53]. In these cases, string models can allow a QCD axion.

Here, we are interested in obtaining the QCD axion also together with a quintaxion in string inspired models [7]. We arrived at this kind of scenario from old solutions of the $\mu$ problem [54] through the composite QCD axion [26]. The superstring axions couples to both the QCD anomaly and the hidden sector anomaly. In this case, we must consider two $\theta$ angles, $\theta_{QCD}$ and $\theta_h$. To have the vacuum at $\theta_{QCD} = \theta_h = 0$, we need two axions. Here, we want to introduce two such axions as the composite one from the hidden-sector squark condensation and a superstring axion. In this case, there exists a problem that an axion with a higher potential corresponds to the smaller axion decay constant and a shallower potential corresponds to the larger decay constant [48]. Since the hidden sector is expected to be at $\Lambda_h \sim 10^{13}$ GeV, the QCD axion generally corresponds to the Planck mass decay constant, grossly violating energy bound [2]. This correspondence is shown in Fig. 5 as real lines. We have to change this behaviour to introduce a reasonable QCD axion.

\[
\begin{align*}
F \bar{F} &\quad \quad \quad \rightarrow \quad \quad \quad F_a = 10^{12} \text{ GeV} \\
F_h \bar{F}_h &\quad \quad \quad \rightarrow \quad \quad \quad F_a \sim M_P
\end{align*}
\]

**Fig. 5.** One naively anticipates that the QCD axion corresponds to a Planck scale decay constant. Our objective is that the QCD axion should corresponds to $F_a \approx 10^{12}$ GeV as shown with the dashed arrows.

It seems that the string axion problem is a real problem. However, if the hidden-sector instanton potential is made shallower than the QCD axion potential, the QCD axion can be saved, as shown in Fig. 5 as dashed lines. In this regards, we adopt the idea of the almost massless h-quarks. But if the h-quark mass is exactly zero, then there is no quintessence. So we must render the h-quarks tiny masses. As we have seen in Sec. 1, the current quark masses come in the determination of the instanton potential. For the hidden-sector vacuum, we
assume almost massless hidden-sector quarks. Also hidden sector gluino mass \((\sim 100 \text{ GeV})\) is much smaller than the hidden sector scale \(\Lambda_h \sim 10^{13} \text{ GeV}\).

A solution of the \(\mu\) problem is assuming a Peccei-Quinn symmetry,\(^4\) forbidding the \(H_1 H_2\) term and \(Q_i \bar{Q}_j\) term in the superpotential. The Peccei-Quinn symmetry, however, allows

\[
W_\mu = \frac{c}{M_p} H_1 H_2 Q_i \bar{Q}_j
\]

(15)

where \(c\) is an order 1 number. If the h-squarks \(Q_i \bar{Q}_j\) condenses at the intermediate scale we obtain a \(\mu\) term at the electroweak scale order. This is one side of a coin. The other side of coin is that if the Higgs fields develops vacuum expectation values at the electroweak scale, \(v \simeq 248 \text{ GeV}\), then the h-quarks obtain tiny masses of order

\[
m_Q \simeq 0.64 \times 10^{-14} \sin 2\beta \ [\text{GeV}]
\]

(16)

where \(\beta = \langle H_2^0 \rangle / \langle H_1^0 \rangle\).

For small instantons, the height of the potential is estimated to be, for \(n\) h-quark flavors and \(N\) number of colors,

\[
\lambda_h^4 \simeq m_Q^4 m_G^N A_h^{4-n-N}
\]

(17)

Thus, the hidden sector potential can be extremely flat, contrary to a naive expectation. It is because the h-quark mass and h-gluino mass are much smaller than the hidden sector scale \(A_h\). Anyway, even though it looks weird, the hidden sector axion potential is much smaller than the QCD axion potential, realizing the dashed line correspondence of Fig. 5 \(^7\). It has been much pronounced by almost the massless quark.

Summarizing, we needed two axions which are provided by (1) superstring axions, either the model-independent axion or model-dependent axions which becomes the quintaxion, providing the dark energy and (2) the axion through introducing a Peccei-Quinn symmetry needed for solving the supergravity \(\mu\) problem. The h-squark condensation introduces another axion which becomes the QCD axion, providing the CDM in the universe.

5 Conclusion

In this talk, I reviewed the current status of the strong CP problem and axion. Regarding the QCD axion, we showed also the current status of the axion search bound, assuming that the QCD axion is the dominant component of the CDM. Furthermore, we speculated that the axionlike particles present in superstring models can be candidates for the source of the dark energy in the universe.

\(^4\) Supersymmetry breaking in supergravity can generate a \(\mu\) term \(^{55}\). But here also, one must assume the absence of \(H_1 H_2\) in one way or another.
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References
1. J. E. Kim, Phys. Rep. 150, 1 (1987); R. D. Peccei, hep-ph/9807514; M. Dine, hep-ph/0011376.
2. J. Preskill, M. B. Wise, and F. Wilczek, Phys. Lett. B120, 127 (1983); L. F. Abbott and P. Sikivie, Phys. Lett. B120, 133 (1983); M. B. Dine and W. Fischler, Phys. Lett. B120, 137 (1983).
3. L. Covi, J. E. Kim, and L. Roszkowski, Phys. Rev. Lett. 82, 4180 (1999); L. Covi, H. B. Kim, J. E. Kim, and L. Roszkowski, JHEP 0105, 033 (2001).
4. J. E. Kim, in Dark matter in astro- and particle physics, Proc. DARK 2002 Conference (Cape Town, South Africa, Feb. 4–9, 2002), ed. H. V. Klapdor-Kleingrothaus and R. D. Viollier, (Springer, Berlin, 2002), p. 246-254.
5. C. G. Callan, R. Dashen, and D. J. Gross, Phys. Lett. B61, 334 (1976); R. Jackiw and C. Rebbi, Phys. Rev. Lett. 37, 172 (1976).
6. M. Chaichian and A. B. Kobakhidze, Phys. Rev. Lett. 87, 171601 (2001).
7. J. E. Kim and H. P. Nilles, Phys. Lett. B553, 1 (2003).
8. T. D. Lee, Phys. Rev. D8, 1226 (1973).
9. M. A. B. Beg and Tsao, Phys. Rev. Lett. 41, 278 (1978); R. N. Mohapatra and G. Senjanovic, Phys. Lett. B79, 283 (1978); H. Georgi, Hadronic J. 1, 155 (1978); G. Segre and H. A. Weldon, Phys. Rev. Lett. 42, 1191 (1979); S. M. Barr and P. Langacker, Phys. Rev. Lett. 42, 1654 (1979).
10. A. Nelson, Phys. Lett. B136, 387 (1984); S. M. Barr, Phys. Rev. Lett. 53, 329 (1984).
11. D. B. Kaplan and A. V. Manohar, Phys. Rev. Lett. 56, 2004 (1986).
12. K. Choi, C. W. Kim, and W. K. Sze, Phys. Rev. Lett. 61, 794 (1988).
13. H. Leutwyler, Nucl. Phys. B337, 108 (1990).
14. K. Choi, Nucl. Phys. B383, 58 (1992).
15. S. Weinberg, Trans. New York Acad. Sci. 38, 185 (1977).
16. G. ’t Hooft, Phys. Rev. D14, 3432 (1976).
17. J. E. Kim, Univ. of Pennsylvania preprint, “Instanton enhancement of $\Delta I = \frac{1}{2}$ nonleptonic interactions”, UPR-0106T (Jan., 1979).
18. A. Irving, C. McNeile, C. Michael, and K. Sharkey, Phys. Lett. B518, 243 (2001); D. R. Nelson, G. T. Fleming, and G. W. Kilcup, Nucl. Phys. Proc. Suppl. 106, 221 (2002) [hep-lat/0110112].
19. I am indebted to Kiwoon Choi for transforming the lattice number to this chiral perturbation parameter.
20. S. Weinberg, Phys. Rev. Lett. 40, 223 (1978); F. Wilczek, Phys. Rev. Lett. 40, 279 (1978).
21. R. D. Peccei and H. R. Quinn, Phys. Rev. Lett. 38, 1440 (1977).
22. E. Witten, Phys. Lett. B149, 351 (1984); Phys. Lett. B155, 151 (1985).
23. K. Choi and J. E. Kim, Phys. Lett. B165, 71 (1985).
24. K. Choi, Phys. Rev. D56, 6588 (1997).
25. J. E. Kim, Phys. Rev. D31, 1733 (1985); K. Choi and J. E. Kim, Phys. Rev. D32, 1828 (1985).
26. E. J. Chun, J. E. Kim, and H. P. Nilles, Nucl. Phys. B370, 105 (1992).
27. J. E. Kim, Phys. Rev. Lett. 43, 103 (1979); M. Shifman, A. Vainshtein, and V. Zakharov, Nucl. Phys. B166, 493 (1980).
28. M. Dine, W. Fischler, M. Srednicki, Phys. Lett. B104, 199 (1981); A. Zhitnitsky, Sov. J. Nucl. Phys. 31, 260 (1980).
29. C. Vafa and E. Witten, Phys. Rev. Lett. 53, 535 (1984).
30. W. A. Bardeen and S. H. H. Tye, Phys. Lett. B76, 580 (1978).
31. H. Georgi and L. Randall, Nucl. Phys. B276, 241 (1986).
32. R. D. Peccei, “A short review of axions”, in Proc. 19th ICHEP (Tokyo, Japan, Aug. 23-30, 1978), ed. S. Homma et. al. (Phys. Soc. Japan, 1979), p. 1045.
33. P. Sikivie, Phys. Rev. Lett. 51, 1415 (1983).
34. N. Iwamoto, Phys. Rev. Lett. 53, 1198 (1984).
35. A. Pantziris and K. Kang, Phys. Rev. D33, 3509 (1986).
36. G. Raffelt and D. Seckel, Phys. Rev. Lett. 60, 1797 (1988); M. S. Turner, Phys. Rev. Lett. 60, 1797 (1988); R. Mayle et. al., Phys. Lett. B203, 188 (1988); T. Hatsuda and M. Yoshimura, Phys. Lett. B203, 469 (1988).
37. K. Choi, K. Kang, and J. E. Kim, Phys. Rev. Lett. 62, 849 (1989).
38. R. L. Davis, Phys. Lett. B180, 225 (1986).
39. P. Sikivie, Phys. Rev. Lett. 48, 1156 (1982).
40. J. Ellis, J. E. Kim, and D. V. Nanopoulos, Phys. Lett. B145, 181 (1984).
41. M. S. Turner, Phys. Rev. D32, 843 (1985).
42. (RB), De Panfilis et.al., Phys. Rev. Lett. 59, 839 (1987); (F), C. Hagmann, P. Sikivie, N. S. Sullivan, and D. B. Tanner, Phys. Rev. D42, 1297 (1990); (LLNL), C. Hagmann et. al., Nucl. Phys. Proc. Suppl. 51B, 209 (1996); (LLNL-MIT), S. J. Azatals et. al., Astrophys. J. 571, L27 (2002).
43. I. Ogawa et. al., Phys. Rev. D53, 1740 (1996); S. Matsuki et. al., Nucl. Phys. Proc. Suppl. 51B, 213 (1996).
44. J. E. Kim, Phys. Rev. D58, 055006 (1998).
45. See, for example, J. E. Kim, Phys. Lett. B564, 35 (2003).
46. S. Perlmutter et. al., Astrophysics. J. 517, 565 (1999); C. L. Bennet et. al., astro-ph/0302207.
47. J. A. Frieman, C. T. Hill, A. Stebins, and I. Waga, Phys. Rev. Lett. 75, 2077 (1995).
48. J. E. Kim, JHEP 9905, 022 (1999); JHEP 0006, 016 (2000); K. Choi, Phys. Rev. D62, 043509 (2000).
49. M. Bronstein, Phys. Z. Sowjetunion 3, 73 (1933); M. Özer and M. O. Taha, Nucl. Phys. B287, 797 (1987); B. Ratra and P. J. E. Peebles, Phys. Rev. D37, 3406 (1988); C. Wetterich, Nucl. Phys. B302, 645 (1988); H. Gies and C. Wetterich, Acta Phys. Slov. 52, 215 (2002); J. A. Frieman, C. T. Hill, and R. Watkins, Phys. Rev. D46, 1226 (1992); R. Caldwell, R. Dave, and P. J. Steinhardt, Phys. Rev. Lett. 80, 1582 (1998); P. Binetruy, Phys. Rev. D60, 063502 (1999); C. Kolda and D. H. Lyth, Phys. Lett. B458, 197 (1999); T. Chiba, Phys. Rev. D60, 083508 (1999); P. Brax and J. Martin, Phys. Lett. B468, 40 (1999); A. Masiero, M. Pietroni, and F. Rosati, Phys. Rev. D61, 023504 (2000); M. C. Bento and O. Bertolami, Gen. Relativ. Gravi. 31, 1461 (1999); F. Perrotta, C. Baccigalupi, and S. Matarrasso, Phys. Rev. D61, 023507 (2000); A. Arbey, J. Lesgourgues, and P. Salati, Phys. Rev. D65, 083514 (2002).
50. K. Choi and J. E. Kim, Phys. Lett. B154, 393 (1985).
51. M. Dine, N. Seiberg, and E. Witten, Nucl. Phys. B289, 585 (1987); J. Atick, L. Dixon, and A. Sen, Nucl. Phys. B292, 109 (1987); M. Dine, I. Ichinose, and N. Seiberg, Nucl. Phys. B293, 253 (1987).
52. J. E. Kim, Phys. Lett. B207, 434 (1988).
53. G. Lazarides and Q. Shafi, Phys. Rev. Lett. 56, 432 (1986); K. S. Babu, I. Gogoladze, and K. Wang, Phys. Lett. B560, 214 (2003); A. G. Dias and V. Pleitez, hep-ph/0308037.
54. J. E. Kim and H. P. Nilles, Phys. Lett. B138, 150 (1984).
55. G. Giudice and A. Masiero, Phys. Lett. B206, 480 (1988).