Quintessence axion revisited in light of swampland conjectures

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Abstract

We point out that the swampland conjectures, forbidding the presence of global symmetries and (meta-)stable de Sitter vacua within quantum gravity, pick up a dynamical axion for the electroweak SU(2) gauge theory as a natural candidate for the quintessence field. The potential energy of the electroweak axion provides an attractive candidate for the dark energy. We discuss constraints from the weak gravity conjecture, from the conjecture of no global symmetry, and from observations, which can be satisfied elegantly in a supersymmetric extension of the standard model.

Keywords: quintessence, axion, swampland conjecture

(Some figures may appear in colour only in the online journal)

Introduction

It has long been a challenging problem to identify any possible consequences of UV quantum gravity for IR physics. For this purpose there has been several conjectures (the so-called swampland conjectures [1–3]) in the literature, which claim necessary conditions for a low-energy effective field theory to have a consistent UV completion within theories of quantum gravity.

One of the most striking conjectures states that a stable de Sitter (dS) vacuum is forbidden in quantum gravity (see [4–8] for recent discussion), as is articulated in the recent de Sitter swampland conjecture [8] and its refinements [9–13]. The conjectures of [8, 12] are motivated by the no-go result of [14], and is better motivated by the swampland distance conjecture

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2, 15] in asymptotic regions of the parameter space where we have parametric control
[12, 16, 17]. While this dS conjecture is still speculative, the conjecture has triggered active
discussions in the community. It was further conjectured that there exists no non-supersym-
metric anti-de Sitter (AdS) vacua [18].

In this Letter, we first point out the inconsistency of the global symmetry with the quantum
gravity makes the θ-angle of the electroweak SU(2) gauge theory a physical parameter. We
then discuss the possibility that the electroweak axion can be regarded as the quintessence
field generating the dark energy which is essential to render the observed dark energy cons-
istent with the conjectural absence of the dS vacua\textsuperscript{5}. We find impressive consistency with
theoretical and observational constraints, if we consider a supersymmetric extension of the
standard model, where a phenomena, the Supersymmetric Miracle, plays crucial roles.

Electroweak axion from swampland conjectures

Let us first present our argument for the electroweak axion. We emphasize again that our argu-
ment is based on a few swampland conjectures.

In the electroweak SU(2) gauge theory, we consider the θ-angle and the associated θ-vac-
umum [19] for each fixed angle θ. One might argue that the θ-angle in SU(2) gauge theory can
be rotated away by anomalies for the global B+L symmetry (here B and L are baryon and
lepton numbers, respectively), and hence is unphysical. However, an exact global symmetry
is forbidden in theories of quantum gravity [20–23], and we expect that the B+L symmetry is
broken by higher-dimensional operators in the standard model, such as the dimension 6 opera-
tor $qqql$ (where $q$ and $l$ are the quarks and the leptons, respectively). In the presence of
B+L breaking, the θ-angle of the electroweak SU(2) gauge theory can no longer be rotated away.

The θ-vacuum is a stable vacuum, and we can consider such vacuum for any value of θ as
a boundary condition of the universe. Here, we assume that the UV quantum gravity provides
a Lagrangian of the effective low energy quantum field theory while it does not control the
boundary condition of the low energy theories. Thus, the vacuum energy is generically of
order of the dynamical scale of the electroweak theory. The existence of such vacua contra-
dicts swampland conjectures, that is, if the vacuum energy is positive this violates the refined
swampland dS conjecture of [12]. The inconsistency with the swampland dS conjecture often
invokes the quintessence field to tilt the scalar potential at the vacuum slightly.

The above arguments point us an interesting possibility of identifying the quintessence
field with a dynamical axion for the electroweak SU(2) gauge theory. Namely the θ-angle is
promoted to the electroweak axion $a$, which couples to the electroweak gauge field as:

$$L_{a\tilde{F}} = \frac{a}{32\pi^2 f_a} \text{Tr} F_{\mu\nu} \tilde{F}^{\mu\nu},$$

where $f_a$ is the decay constant. Here, it should be emphasized that we make a very strong
assumption that the electroweak axion does not have any other interaction terms beyond the
one in equation (1). If there are other shift symmetry breaking terms, the axion potential may
have many local minima with positive vacuum energy in general, which are also excluded by
the swampland conjectures. Besides, the additional contributions to the scalar potential of the
electroweak axion can spoil the numerical success of the size of the dark energy generated by
the electroweak instanton which we will see shortly.

\textsuperscript{5} A similar discussion for the QCD axion [57–59] can be found in the recent paper [60].
Incidentally, the presence of two light axions (the electroweak axion and the QCD axion) has unexpected virtue, in that the two light axions removes the tension between Majorana neutrinos and the conjectured absence of the non-supersymmetric stable AdS vacua, see section 5.2.2 of [24].

Electroweak axion as quintessence

The shift symmetry $a \rightarrow a + (\text{const.})$ of the axion is broken in the presence of the $qqql$ operator, and non-perturbative instanton effects generate an axion potential [25],

$$V(a) = \Lambda_a^4 \left( 1 - \cos \frac{a}{f_a} \right),$$

in the one-instanton approximation. Here $\Lambda_a$ is the dynamically-generated scale, whose naive estimate is given by

$$\Lambda_a^4 = M_{Pl}^4 e^{-\frac{16\pi}{\alpha_2(M_{Pl})}} \simeq 10^{-130} M_{Pl}^4 \ll M_{Pl}^4,$$

where we use the value of the electroweak coupling constant $\alpha_2 = g_2^2/(4\pi)$ at the Planck scale $\alpha_2(M_{Pl}) = 1/48$, and $M_{Pl} \simeq 2 \times 10^{18}\text{GeV}$ is the reduced Planck scale. The axion potential is dominated by small-size instanton contributions since it requires insertions of the higher-dimensional $qqql$ operators.

While the potential (2) is problematic [11] in the original version [8] of the de Sitter swampland conjecture (see also [26–29]), the refinement [12] is compatible with the axion potential as it is vanishing at the minimum at $a = 0$ (mod $2\pi f_a$). In this Letter we assume the refined dS conjecture of [12].

The curious observation is that the value of $\Lambda_a$ is close to the current energy scale of the dark energy [30, 31],

$$\Lambda_0^4 \simeq 10^{-120} M_{Pl}^4 \simeq O(\text{meV})^4.$$  

This estimate shows that the identification of the idea of the quintessence axion [32–34] is highly successful, where the dark energy can be explained by the dynamical electroweak axion. Here the shift symmetry of the axion ensures the flatness of the quintessence potential against possible quantum corrections.

While the conjectural absence of the de Sitter vacua has already been used as motivations for quintessence [35–37] explanation of dark energy (see [8, 38–41] for recent discussion), we have shown here that the conjecture brings a natural candidate for the quintessence, the quintessence axion.

Constraints from the weak gravity conjecture

To this point, we have motivated the electroweak quintessence axion [30, 31] from the one of the swampland conjectures, namely the absence of the dS and AdS vacua. It turns out, however, this scenario is some tension with another swampland conjecture, the weak gravity conjecture [42].

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6 The axion potential may take a negative value at the minimum which can be consistent with the swampland conjectures on stable AdS vacua as long as we assume that they are unstable quantum mechanically.
The weak gravity conjecture sets an upper bound on the decay constant

$$f_a \lesssim \frac{M_{\text{Pl}}}{S_{\text{inst}}}, \quad (5)$$

where $S_{\text{inst}} \simeq 2\pi/\alpha_2(M_{\text{Pl}})$ is the size of instanton action. Notice we have chosen the value of $\alpha_2$ here to be evaluated at the UV scale and not at the IR scale, since as we have described above the $\theta$-angle is physical only after taking into account higher-dimensional $(B+L)$-breaking operators.

Since $S_{\text{inst}} \sim O(100)$ for the electroweak instanton, we obtain the bound

$$f_a \lesssim O(10^{16} \text{GeV}). \quad (6)$$

This is parametrically smaller than the value $f_a \sim M_{\text{Pl}} \sim 10^{18} \text{GeV}$, which was assumed in [30].

While the energy scale of dark energy (3) is not affected by the value of the decay constant $f_a$, such a low value of $f_a$ predicts a larger value for the mass $m_a$ for the quintessence axion [30]:

$$m_a^2 \simeq \frac{\Lambda^4}{f_a^2} \simeq \frac{H_0^2 M_{\text{Pl}}^2}{f_a^2} \gtrsim H_0^2 = (2 \times 10^{-33} \text{eV})^2, \quad (7)$$

where $H_0$ is the present-day value of the Hubble constant. This means that the quintessence field has already started rolling down the potential, and the slow-roll condition $M_{\text{Pl}}|V'(a)| \ll V(a)$ is no longer satisfied. This is problematic for the axion as a dark energy candidate.

One possibility to circumnavigate this problem is to fine-tune the initial condition of the quintessence to be close to the local maximum $a \sim \pi f_a$, so that the quintessence, while it has already started rolling down the potential, is still located close to the local maximum.

This hilltop quintessence scenario [43], however, has a severe fine-tuning problem in our context. Indeed, if we start with a small initial value of the displacement $\delta a = a - \pi f_a$, we expect that $\delta a$ will grow. We can estimate this growth from the equation of motion for the axion

$$\ddot{a} + 3H_0 \dot{a} = -V'(a), \quad (8)$$

where $H_0$ is the Hubble constant. Here we use the present-day value of the Hubble constant since we are interested in the era when the dark energy dominates the potential energy of the Universe. When the displacement $\delta a$ is small, we can approximate equation (8) as

$$\dot{\delta a} + 3H_0 \delta a = \frac{3H_0^2 M_{\text{Pl}}^2}{f_a^2} \delta a, \quad (9)$$

where we used $\Lambda^4 \simeq 3H_0^2 M_{\text{Pl}}^2$. Solving this equation we estimate that $\delta a$ grows exponentially $(\delta a(t) \propto \exp(\sqrt{3H_0 M_{\text{Pl}}/f_a} t))$, so that after the axion starts rolling the total growth until today ($t \sim O(H_0^{-1})$) is of order $\exp(O(M_{\text{Pl}}/f_a)) = \exp(O(100))$. In order to ensure slow-roll condition until now, initial displacement when the axion starts rolling needs to be constrained in a tiny window

$$\left| \frac{\delta a}{f_a} \right| \ll \frac{f_a}{M_{\text{Pl}}} e^{-O(100)}. \quad (10)$$

\[7\] See also [34]. Similar exponential fine-tuning was discussed recently in [61] for curvatons.
One moreover expects that such an extreme fine-tuning is incompatible with the fluctuations generated during the inflationary era.

We therefore conclude that the hilltop quintessence scenario, as required by the small value $f_a \lesssim O(10^{16} \text{ GeV})$ of the decay constant (6), is strongly disfavored.

### Inclusion of heavy particles

Despite the difficulties mentioned above, we can save the electroweak quintessence axion scenario by taking into account the effects of heavy particles.

Recall that in the weak gravity conjecture bound (5) we have substituted the value $S_{\text{inst}} \sim O(100)$ which was obtained by renormalization group (RG) running in the SU(2) gauge group all the way up to the Planck scale, assuming that there are no massive particles modifying the RG running of $\alpha_2$. However, there is no strong justification for this assumption.

Suppose that the RG running of $\alpha_2$ is changed so that we have a larger value of $S_{\text{inst}}$ at the Planck scale. For example, if $S_{\text{inst}} = O(1)$ we have $f \sim M_{\text{Pl}}$. The axion mass (7) can then be taken to be $m_a \sim H_0$, which makes the fine-tuning condition on the field value in equation (10) much less severe to achieve the slow-roll conditions around the hill-top region. It should be also noted that the value $f_a \lesssim M_{\text{Pl}}$ of the decay constant is also compatible with the refined swampland distance conjecture [2, 15], which restricts the field range for the axion to be $\Delta a \lesssim O(M_{\text{Pl}})$.

There is, however, still one problem if one changes the value of the coupling constant at the Planck scale $\alpha_2(M_{\text{Pl}})$—the same coupling constant appears in the dynamical energy scale $\Lambda_a$ of the quintessence (3), and once we modify the value of $\alpha_2(M_{\text{Pl}})$ we are spoiling the crucial observation $\Lambda_a \simeq \Lambda_0$, which was the very starting point for our quintessence axion scenario.

Interestingly, this problem is solved elegantly in supersymmetric theories, which we turn next.

### Supersymmetric miracle

Let us consider the minimal supersymmetric standard model (MSSM). The dynamical scale of the axion potential is computed in the instanton calculus as [30]

$$
\Lambda_a^4 \simeq e^{-\frac{2\pi}{\alpha_2(M_{\text{Pl}})_{\text{MSSM}}}} e^{10 m_{\text{SUSY}}^3 M_{\text{Pl}}}.
$$

(11)

Here $\alpha_2(M_{\text{Pl}})_{\text{MSSM}}$ is the SU(2) gauge coupling constant at the Planck scale for the MSSM. The mass scale $m_{\text{SUSY}} \simeq O(\text{TeV})$ is the scale for spontaneous supersymmetry breaking, whose exact value is not too important for what follows.

In the MSSM, the $B+L$ symmetry is broken by Planck-suppressed dimension 5 operators $QQQL$ [44, 45]. These operators induce too rapid proton decay for $m_{\text{SUSY}} \simeq O(\text{TeV})$. To suppress the dimension 5 operators, we assume the Froggatt–Nielsen (FN) symmetry [46], following [30]. For now, we assume the FN symmetry is a global symmetry. We also assume

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8 For $f_a = O(M_{\text{Pl}})$, the axion potential can be modulated by the higher harmonics or gravitational instanton effects. Since the domain of the axion field is compact, the axion potential have local maximum even with those effects, and hence, the slow-roll conditions can be satisfied around the hill-top region without severe fine-tuning as long as the overall scale of the axion potential is not altered significantly.

9 The deviation from $\Lambda_a \simeq \Lambda_0$ might open up another interesting application of the EW axion such as a candidate for the ultralight (fuzzy) matter [62].

10 Another possibility is to gauge a discrete subgroup of the FN symmetry. For example, when the FN charges are chose as in table 2 of [30] the $Z_{10}$ subgroup is anomaly free and can be gauged.
the global $R$-symmetry given in [30]. The FN symmetry is spontaneously broken by $\epsilon$, where the value of $\epsilon$ is taken to be $\epsilon \approx 1/17$ [47] to explain the quark/lepton masses and mixing angles. With the charge assignment in table 2 of [30], we confirm that the proton decay via the $QQQL$ operators are sufficiently suppressed.

If we do not include any heavy particles beyond the MSSM we can substitute the value $\alpha_2(M_{Pl})|_{MSSM} \approx 2/11$, so that we obtain

$$\Lambda_4^a \approx \left( \frac{\epsilon}{1/17} \right)^{10} (\text{meV})^4,$$

which nicely reproduces the scale $\Lambda_0$ of the dark energy [30].

Suppose that we include a pair of heavy massive particle $X, \bar{X}$ in some representation $R$ of SU(2) gauge group. When the mass of $X, \bar{X}$ is at an intermediate energy scale $M_X$, then the RG running of the coupling constant in the one-loop approximation is modified as

$$\alpha_2^{-1}(M_{Pl})|_{XX} = \alpha_2^{-1}(M_{Pl})|_{MSSM} + \frac{2T_R}{2\pi} \log \frac{M_X}{M_{Pl}},$$

where $T_R$ is the Dynkin index of the representation $R$. Such heavy particles also change the zero modes, so that we need to insert operators $M_X XX$ to cancel the instanton zero modes, leading to the factor $(M_{Pl}/M_X)^{2T_R}$. Interestingly, these two effects cancel with each other for a supersymmetric theory, leaving the dynamical scale $\Lambda_4^a$ invariant:

$$\Lambda_4^a|_{XX} \simeq e^{-\frac{2\pi}{\alpha_2(M_{Pl})|_{MSSM}} \left( \frac{M_X}{M_{Pl}} \right)^{2T_R} \epsilon^{10} m_{\text{SUSY}}^2 M_{Pl}}$$

$$= e^{-\frac{2\pi}{\alpha_2(M_{Pl})|_{MSSM}} \epsilon^{10} m_{\text{SUSY}}^2 M_{Pl}} \simeq \Lambda_4^a|_{MSSM},$$

which nicely reproduces the scale $\Lambda_0$ of the dark energy [30].

![Figure 1](image-url). The one-loop RG running of the SU(3) $\times$ SU(2) $\times$ U(1) gauge coupling constants as a function of the energy scale $\mu$. We color $\alpha_1^{-1}, \alpha_2^{-1}, \alpha_3^{-1}$ by purple, light blue, dark blue, respectively. We can achieve both (1) $\alpha_3(M_{Pl}) \approx 2\pi$ and (2) gauge coupling unification at the Planck scale $M_{Pl}$ by including an SU(2) triplet and an SU(3) octet at $10^{12}$ GeV, together with 4 pairs of GUT-like multiplets $\mathbf{5}$ and $\bar{\mathbf{5}}$ at $\sim 10$ TeV. The RG running in this case, with the supersymmetry breaking scale $m_{\text{SUSY}} \approx 1$ TeV, is shown as solid lines. This is contrasted with the case of the MSSM, shown as dotted lines.
We call this cancellation ‘supersymmetric miracle’. Due to the Supersymmetric Miracle, $\alpha^2(M_{Pl})$ can be achieved while the relation $\Lambda_a \approx \Lambda_0$ intact. While the one-instanton approximation is not reliable for a large value of the coupling constant $\alpha^2(M_{Pl}) \sim 2\pi$, we expect that the estimation of the energy scale $\Lambda_a$ (12) will not be significantly affected by this subtlety.

Once $\alpha^2(M_{Pl}) \sim 2\pi$ is achieved, the size of the instanton action is of $S_{\text{inst}} = O(1)$, which makes the large decay constant, $f_a \sim M_{Pl}$, marginally consistent with the weak gravity conjecture in equation (5). Consequently, the axion mass $m_a$

$$m_a^2 \approx \frac{\Lambda_a^4}{f_a^2} \approx \left(\frac{\epsilon}{1/17}\right)^{10} \left(\frac{2 \times 10^{18}\text{GeV}}{f_a}\right)^2 (7 \times 10^{-34}\text{eV})^2, \quad (15)$$

is achieved. In this way, the electroweak action serves as the slow-rolling quintessence field consistently with the weak gravity conjecture.

There are many scenarios for realizing $\alpha^2(M_{Pl}) \sim 2\pi$ (so that $f \sim M_{Pl}$). For example, we can simply include 3 SU(2) triplets at $\sim 10^7$ GeV. As another possibility consistent with a coupling unification as in grand unified theories (GUT), let us consider an SU(2) triplet and an SU(3) octet at $10^{12}$ GeV. We see that all gauge couplings meet at the Planck scale [48, 49]. The value $\alpha^2(M_{Pl}) \sim 2\pi$ can then be achieved by including 4 pairs of GUT-like multiplets $5$ and $\bar{5}$ at $\sim 10$ TeV (see figure 1). In any of those scenarios, $\Lambda_a \sim \Lambda_0$ is not affected by the Supersymmetric Miracle.

Incidentally, in a supersymmetric theory the axion is complexified, so that we have another scalar (the saxion)\(^{11}\). The saxion causes a cosmological problem similar to the Polonyi problem [50–54]. This problem can easily be solved if the saxion strongly couples to the inflaton [55, 56].

Finally, let us discuss the issues of the global FN symmetry and the $R$-symmetry. As the swampland conjecture precludes the existence of exact global symmetries, those symmetries cannot be exact. On the other hand, the sizes and the patterns of the explicit breaking due to the quantum gravity are under debate. Thus, for example, if the sizes of the explicit breaking are of $O(10^{-1})$ or smaller for unit charges, the successful features of the EW axion quintessence are not spoiled\(^{12}\).

In summary, in this Letter we pointed out that the swampland conjectures bring a dynamical axion for the electroweak SU(2) gauge theory as an interesting candidate for the quintessence field. We also discussed constraints both from the weak gravity conjecture, from the conjecture of no global symmetry, and from observations. We found that those constraints are satisfied elegantly in a supersymmetric extension of the standard model. There, the large decay constant, $f_a \sim M_{Pl}$, is achieved without spoiling the successful relation, $\Lambda_a \sim \Lambda_0$. It is a fascinating question to explore if our scenario can really be realized in a specific setup inside theories of quantum gravity, such as string theory.

\(^{11}\)The existence of the saxion is predicted more generally from one of the swampland conjectures in [2].

\(^{12}\)Without any global symmetry, it is not possible to suppress a local operator which can absorb all the fermion zero modes associated with the electroweak instanton. Such an operator can change the estimation of $\Lambda_a$ in equation (11) drastically. Once we allow the global $R$-symmetry, on the other hand, such a local operator is highly suppressed even if it is an approximate symmetry.
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