Bosonic Coherent Motions in the Universe

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We review the role of fundamental spin-0 bosons as bosonic coherent motion (BCM) in the Universe. The fundamental spin-0 bosons have the potential to account for the baryon number generation, cold dark matter (CDM) via BCM, inflation, and dark energy. Among these, we pay particular attention to the CDM possibility because it can be experimentally tested with the current experimental techniques. We also comment on the panoply of the other roles of spin-0 bosons—such as those for cosmic accelerations at early and late times.

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I. INTRODUCTION

Recent cosmological observations [1, 2] confirm the eight decade old Zwicky’s proposal [3] that the Universe contains a large amount of dark matter (DM). The DM profile has been measured accurately enough to pin-point DM to “it is cold dark matter (CDM) [1].” The bosonic coherent motion (BCM) can be CDM [4–6] if the coherent-boson lifetime is long enough to have survived until now [7–10]. The axion proposed to solve the strong CP problem [11–14] is fitting to this BCM scenario [15]. The BCM is one of many possibilities of CDM scenarios [15] which accounts for only 27% in the energy pie of the Universe. The dominant portion, 68%, in the energy pie is the homogeneous energy density, at least up to the $10^3$ Mpc scale, which is usually referred to dark energy. Dark energy (DE) being homogeneous cannot be accounted for by corpuscular particles but may be accountable by the cosmological constant or by some vacuum expectation value (VEV) of spin-0 boson(s) [16]. The visible particles (mostly atoms in the energy count) constitute only 5% in the energy pie.

If we accept the Big Bang cosmology from the earliest possible time, $10^{-43}$ s, the success of the Standard Model of particle physics is based on the assumption of very tiny DE of order less than $10^{-46}$ GeV$^4$ because the age of the Universe is very long $\simeq 13.8$ Gy [1, 2]. So, the DE problem or the theoretical cosmological constant problem [24] is not only the problem in cosmology but also a problem in particle physics. Out of despair, many adopt the anthropic scenario for the cosmological constant problem [24, 26]. For the anthropic solution to work, the cosmological constant must be a free undetermined parameter in particle physics, as integration constants of Refs. [27–29]. In a deterministic theory such as string theory, possible cosmological constants must be allowed near 0 for our Universe to have adopted one of these, which is the reason trying to have as many as $10^{120}$ models, to pack the vacua with separation between them by $(10^{-3}$ eV)$^4$, from string theory [26]. But, all those $10^{120}$ vacua must allow three family SMs, and satisfy the known SM phenomena such as the GUT scale weak mixing angle $\sin^2 \theta_W = 3/8$ [30, 32], etc. But, we have only a handful of minimal supersymmetric Standard Models from string theory satisfying the requirements [33, 37]. Or, a Standard Model solution with DE $\simeq 10^{-47}$ GeV$^4$ has to be found so that the anthropic argument chooses it. This search seems more difficult than finding a vanishing cosmological constant solution theoretically. At present, we can say that the anthropic solution in string theory has not worked out yet. Therefore, in the Standard Model and in its supergravity extension, it is fair to say that the cosmological constant is assumed to be zero.

By observing the luminosities of Type-I supernovae [38, 39], the recent acceleration of the Universe has been established. So, explaining the DE scale of $10^{-47}$ GeV$^4$ -- $10^{-46}$ GeV$^4$ became an important topic [14, 21]. In Table I, we list several ideas proposed to account for this recent acceleration of the Universe. Both the high scale inflation [44, 47] and the recent acceleration [38, 39] in the Universe are based on the assumption of vanishing cosmological constant.
| Ideas                        | Description [scalar S or pseudoscalar P] | Disc. sym. Fine-tune From string |
|-----------------------------|------------------------------------------|----------------------------------|
| MOND                       | Change Newtonian gravity. [No boson]     | Irrelevant Yes Not yet           |
| Anthropic principle        | Out of many possible vacua, only those  | Irrelevant Irrelevant Not yetb  |
| Quintessence               | With a runaway $V \propto 1/\phi^n$ ($n > 0$). [S] | No Yesc Not yet                |
| Dilaton                    | P-Gold. boson from dilaton sym. [S]     | No Yesd Not yet                 |
| U(1)$_{DE}$ Goldstone      | P-Gold. boson from U(1)$_{DE}$ sym. [P]  | Yes No Yeso                    |

TABLE I: Typical DE models with a few pseudo-Goldstone bosons originating from global symmetries. a Refs. 40, 41. b Ref. 26. c Ref. 42. d Ref. 43. e Ref. 44.

To determine the VEV of a scalar field, say $\phi$, one must consider all the allowed effective terms at low energy. At each interaction point, suitable symmetry requirements must be satisfied. A typical mass scale of $\phi$ is given by the effective mass term $m^2 |\phi|^2$. In Fig. 1, we consider only two diagrams with the dimension 4 ($d = 4$) couplings. If each $d = 4$ vertex of Fig. 1 satisfies the global phase symmetry, the two-loop and one-loop mass terms do not break the global symmetry. On the other hand, each $d = 4$ vertex satisfies the dilaton symmetry (requiring just $d = 4$ couplings) but the diagrams of Fig. 1 are $d = 2$ terms which of course break the dilaton symmetry. One well-known model breaking the dilaton symmetry at the one-loop quantum level, including the $d \geq 6$ terms, is the Coleman-Weinberg model 48. Therefore, it is not likely that a consistent calculation of a small DE scale can be performed by introducing the dilaton symmetry. However, some global phase symmetry may be suitable for this.

In Sec. II we present the focus points of this review: the BCM scenarios and the axion detection experiments. In Sec. III we point out the difficulty of obtaining zero cosmological constant theoretically. In Sec. IV we mini-

review the inflationary cosmology, in particular in view of the recent BICEP2 data. In Sec. V we discuss the subject of this review: why the role of fundamental spin-0 particles are important in cosmology.

II. SPIN-0 BOSON FILLING THE UNIVERSE

After the discovery of a fundamental spin-0 scalar particle (the Brout-Englert-Higgs boson) at the LHC, it is timely to study the roles of fundamental spin-0 bosons in the Universe. It is very interesting to note that fundamental spin-0 bosons have been employed to account for the mothers of atoms (i.e. baryon number generation via the Affleck-Dine mechanism 49), CDM via BCM 4, DE via a transient cosmological constant 50-54, and even the vacuum energy needed for the high scale inflation 45-47. Among these, we focus on CDM via BCM in this review because similar ideas can be applicable to DE and inflation models. Another attractive point discussing CDM via BCM is that it can be experimentally proved in the near future 55.

FIG. 1: Diagrams leading to dimension 2 interactions with dimension 4 coupling at each vertex.

We are familiar with the ether idea of the late 19th Century, filling out the Universe. The VEV idea of spin-0 particles used for breaking global symmetries 56 and gauge symmetries 57 is a kind of ether. If a scalar field $\phi$ has a universal value over the entire Universe, any operation of the type ‘Poincare transformation’ does not notice a change. Thus, the VEV of a scalar field, $\langle \phi \rangle$, respects the Poincare symmetry. But, if $\phi$ is a complex field, then the VEV breaks the phase transformation symmetry, i.e. breaks a global phase symmetry 58. Even though the Brout-Englert-Higgs mechanism 57,54 for breaking gauge symmetries is not a monopoly of spin-0 particles 60,61, now the role of spin-0 particles becomes more important, especially after a hint of large tensor-to-scalar ratio $r$, based on the BICEP2 observation 62.

Let us denote scalar and pseudoscalar particles as $s$ and $a$, respectively. Scalar particles transform under the parity operation as $P: s(x) \rightarrow +s(-x)$, and pseudoscalar particles transform as $P: a(x) \rightarrow -a(-x)$. If they are components of a complex field, it is usually represented as the radial and phase fields, respectively, $\phi = s e^{ia/f}$ where $f$ is a mass parameter. Thus, the complex field
transforms under parity as $P : \phi(x) \rightarrow \phi^*(-x)$. Any pseudoscalar field represented as a phase can be represented by an angle field with the angle defined in the range $[0, 2N\pi)$, where $N$ is the domain-wall number. A Goldstone boson arising from breaking a global phase symmetry by the VEV $v$ is a pseudoscalar field $a$ defined as

$$\langle \phi \rangle = \frac{v + s}{\sqrt{2}} e^{in/2}, \quad \langle s \rangle = 0, \quad (a) = [0, 2N\pi f).$$ (1)

A. Cosmology with BCM

On the flat Friedmann-Lemaître-Robertson-Walker cosmological background space described by the line element $ds^2 = -dt^2 + a^2(t)\delta_{ij}dx^i dx^j$, the evolution of the classical scalar field $\phi$, (i.e., the evolution of the VEV of $\phi$), is given by

$$\frac{d^2}{dt^2} \langle \phi \rangle + 3H \frac{d}{dt} \langle \phi \rangle + V'(\langle \phi \rangle) = 0,$$ (2)

where $H = \dot{a}/a$ is the Hubble parameter and $V' = (d^2/d\phi^2) V$ is a derivative of the potential $V$ (a dot represents a derivative with respect to the cosmic time $t$). With a discrete symmetry $\phi \rightarrow -\phi$, the leading term of $V'$ is the mass term $m^2 \langle \phi \rangle$. When $\langle \phi \rangle$ moves very slowly, we can neglect the second derivative term, and the evolution equation gives $3H \dot{\phi} \simeq -m^2 \phi$. $\langle \phi \rangle$ starts to change rapidly when $H$ becomes small enough to satisfy $3H \simeq m$. After this condition is met, $\langle \phi \rangle$ oscillates rapidly, as shown in Fig. 2, which is interpreted as the BCM of $\phi$.

As mentioned above, the VEV $\langle \phi \rangle$ is assumed to be the same over the whole Universe for the Poincaré invariance, otherwise the invariance is broken. In the Universe, this homogeneity is subtly broken. The inflation manages different scales of quantum fluctuations enter the horizon at different scales, basically breaking the homogeneity. A given scale condenses gravitationally. The VEV in that scale evolves according to Eq. (2), and describes the BCM of $\langle \phi \rangle$.

From the Friedmann equation we have $3H^2 M_P^2 = \rho$, where $\rho$ is the energy density of the Universe and $M_P$ is the reduced Planck mass ($M_P = 2.4 \times 10^{18}$ GeV). Denoting the time at the onset of oscillations of $\langle \phi \rangle$ as $t_1$, the condition for determining $t_1$ is

$$\sqrt{\frac{3\rho(t_1)}{M_P^2}} = m(t_1).$$ (3)

These oscillations are equivalent to a gas of $\phi$ particles of low-momentum. This kind of spin-0 particle coherent motion was first discussed in Ref. [4] for the case $\phi$ = axion. It is known that the BCM behaves like CDM because of the low-momentum. Thus, the number and energy densities are given by

$$n = m \langle \phi \rangle^2, \quad \rho = m^2 \langle \phi \rangle^2.$$ (4)

We denote the current age of the Universe as $t_U$. Depending on $t_1$ and $t_U$, we can classify BCMs as

- **BCM**: If $t_1 < t_U$, the currently oscillating vacuum $\langle \phi \rangle$ is BCM. The BCM can be classified into the following two sub-categories.

  - **BCM1**: The lifetime of $\phi$ is long enough, $\tau_\phi > t_U$. Then, the oscillating BCM contributes to the CDM amount. The QCD axion belongs here.
  
  - **BCM2**: The lifetime of $\phi$ is short, $\tau_\phi < t_U$. Then, all $\phi$ quanta decay already, producing SM particles. The inflaton with $\tau_\phi \sim 10^{-36}\text{s}$ belongs here and reheating after inflation gives the beginning of the radiation-dominated Universe.

- **CCtmp**: Temporary cosmological constant. On the other hand, if $\langle \phi \rangle$ has not oscillated yet, then $t_1 > t_U$ and $\langle \phi \rangle$ stays there now, behaves like a cosmological constant, but it is a temporary phenomenon and will eventually become BCM1 after $t_1$. For this to be satisfied, the mass is around $10^{-38}\text{eV}$ with a trans-Planckian decay constant $\delta_{\phi}$.

  If $V(\langle \phi \rangle)$ describes **CCtmp**, the equation of state $w_\phi$, characterized by the field energy density $\rho_\phi = \frac{1}{2} \phi^2 + V(\phi)$ and the pressure $P_\phi = \frac{1}{2} \phi^2 - V(\phi)$, is a useful parameter,

$$w_\phi \equiv \frac{P_\phi}{\rho_\phi} = \frac{\frac{1}{2} \phi^2 - V(\phi)}{\frac{1}{2} \phi^2 + V(\phi)}.$$ (5)

Provided that $\frac{1}{2} \phi^2 \ll V(\phi)$, $w_\phi$ is close to $-1$, behaving like the cosmological constant. In order to realize the recent acceleration, we require the condition $w_\phi < -\frac{1}{3}$.

B. Scalar particles

The Brout-Englert-Higgs boson is the only known fundamental scalar field. The other scalar most widely used in particle theory is dilaton, the scalar Goldstone boson arising from breaking the dilatonic symmetry. The effect of dilatonic symmetry on the cosmological constant
problem has been discussed extensively in Ref. [43]. For the solution, however, a fine-tuning is needed. The obvious effect of a VEV of a scalar field $s$ in cosmology is the universal constant on the right-hand-side of the Einstein equation.

The scalar-field cosmology in the presence of a barotropic perfect fluid was studied in 1980–90s [43, 64–72], even before the discovery of the recent cosmic acceleration. This was chiefly motivated by the “missing matter problem” in 1980s. In 1990, Fukugita et al. [73] tested cosmological models against observations of the number count of faint galaxies and showed that these data favor the Universe with low matter density (i.e., matter is missing). In the abstract of their paper they stated that “Furthermore, it is shown that the best agreement with the data is obtained with a sizable cosmological constant, including the case of zero curvature model as predicted by inflation.” In addition, it was already known in the early 1990s that the presence of a cosmological constant can make the age of the Universe longer such that it is consistent with the age of oldest globular clusters.

If the cosmological constant originates from a vacuum energy appearing in particle physics, it is vastly larger than the today’s average cosmological density [28]. Because of this problem, people tried to construct dynamical cosmological constant models in which the energy density of cosmological constant varies in time, basically belonging to a kind of CCtmp. For example, if we consider a dilaton field $\phi$, the cosmological constant depends on $\phi$ by transforming the dilatonic action to the so-called Einstein-frame action (in which the dilaton does not have a direct coupling with the Ricci scalar) [43, 65].

Exponential potentials often arise from the curvature of internal spaces associated with the geometry of extra dimensions (so called “modulus” fields) [74, 77]. Inspired by this, the exponential potential $V(s) = V_0 e^{-\lambda s M_P}$ has been used, with constant parameters $V_0$ and $\lambda$. There are two distinct fixed points on the flat Friedmann-Lemaitre-Robertson-Walker cosmological background space [19, 72]: (a) the scaling solution, and (b) the scalar-field dominated solution.

For $\lambda^2 > 3(1 + w_m)$, where $w_m$ is the equation of state for a background fluid, the solutions approach the scaling fixed point (a), characterized by the field density parameter $\Omega_s = 3(1 + w_m)/\lambda^2$ and the field equation of state $w_s = w_m$ [67, 71, 72]. Even for the initial conditions where $\rho_s$ is larger than $\rho_m$ in the early radiation-dominated era, the field eventually enters the scaling regime in which $\rho_s$ is proportional to $\rho_m$ with $\rho_s/\rho_m = constant < 1$. The field energy density of the scaling solution contributes to the total energy density of the Universe, but it does not lead to the cosmic acceleration. For $\lambda^2 < 3(1 + w_m)$, there exists the scalar-field dominated fixed point (b), characterized by $\Omega_s = 1$ and $w_s = -1 + \lambda^2/3$. The late-time cosmic acceleration can be realized for $\lambda^2 < 2$. Since in this case the point (b) is also stable, the scalar field can be the source of DE. For $\lambda^2 < 2$, the scalar potential is quite shallow, so the field density in the early Universe needs to be much smaller than the background energy density (unlike the scaling solution discussed above).

After the discovery of the recent cosmic acceleration in 1998, the cosmological dynamics of “quintessence” (a canonical scalar field responsible for DE) were studied in detail for several different potentials [76, 77]. One example is the inverse power-law potential $V(s) = M^{4/n} s^{-n}$, where $M$ and $n$ are positive constants. This potential can arise in globally supersymmetric QCD theories [78]. The Universe enters the stage of cosmic acceleration for the field value larger than $s_0 \approx M_P$. Since $V(s_0)$ is of the order of $H_0^2 M_P^2$, one can estimate the mass scale $M$ as $M \approx 10^{(46−19n)/(4+n)}$ GeV. For $n = O(1)$, this energy scale can be compatible with that appearing in particle physics.

In the presence of a perfect fluid with the equation of state $w_m$, there exists a so-called tracker solution for the potential $V(s) = M^{4/n} s^{-n}$. The tracker is characterized by a common, cosmic evolutionary trajectory that attracts solutions with a wide range of initial conditions [77]. The field equation of state along the tracker is given by $w_s = (w_m n - 2)/(n + 2)$, which corresponds to $w_s = -2/(n + 2) > -1$ during the matter era. The slope of the potential $\lambda = -M_P V_s/V = n M_P/s$ gets smaller with the growth of $s$, so $w_s$ approaches $-1$ in the future. The inverse power-law potential belongs to a class of freezing quintessence models [73] in which the evolution of the field gradually slows down.

There is another class of quintessence models, dubbed thawing models [79], in which the field has been frozen by Hubble friction and then it starts to evolve after the Hubble parameter drops below the field mass $m$. In this case the field equation of state $w_s$ is close to $-1$ at the initial stage, but it starts to grow at the late cosmological epoch. The field mass $m_s$ responsible for dark energy corresponds to $m_s \approx 10^{-33}$ eV [63]. The representative potential of thawing models is that of a pseudo-scalar field arising from breaking the global $U(1)$ symmetry (which we will explain more details in Sec. II C).

If we consider a scalar field $\phi$ non-minimally coupled to the Ricci scalar $R$ (like dilaton), this gives rise to a coupling with non-relativistic matter in the Einstein frame [80]. The fifth force induced by such a matter coupling needs to be suppressed in the solar system. There are several ways to suppress the propagation of the fifth force in local regions of the Universe.

One is the so-called chameleon mechanism [81], under which the mass of a scalar degree of freedom is different depending on the matter densities in the surrounding environment. If the effective mass is sufficiently large in the regions of high density, the coupling between the field and non-relativistic matter can be suppressed by having a thin shell inside a spherically symmetric body.

1 However, the scalar in this case is composite.
In Brans-Dicke theory (including $f(R)$ gravity) \[82\] it is possible to suppress the propagation of the fifth force by designing the field potential $V(\phi)$ appropriately \[83,88\]. Another is the so-called Vainshtein mechanism \[89\], under which nonlinear scalar-field self interactions can suppress the fifth force at short distances even in the absence of the field potential. The self interactions of the form $(\partial \phi)^2 \phi$, which correspond to the Lagrangian of covariant Galileons \[90\], can lead to the decoupling of the field $\phi$ from matter within a radius much larger than the solar-system scale \[91,92\].

C. Pseudoscalar particles

Most pseudoscalar particles observed so far are pseudo-Goldstone bosons. Let $a, \Lambda$ and $f$, respectively, be a Goldstone boson from a spontaneously-broken global $U(1)$ symmetry, the dominant explicit symmetry breaking mass parameter, and the decay constant. Then, the mass of $a$ is

$$m_a = c_a \frac{\Lambda^2}{f},$$

where $c_a$ is the number given by the explicit symmetry breaking terms. For the QCD axion, the breaking of the $U(1)$ symmetry is given by the QCD anomaly and we have $c_a \Lambda^2 = [Z^{1/2}/(1 + Z)] f_\pi \Lambda \Pi$, with $Z = m_u/m_d$ where $f_\pi, m_\pi, m_u, m_d$ are neutral-pion decay constant, its mass, and $u$ and $d$ quark masses \[93\]. If the explicit breaking term is given by $V_{\text{br}} = -(\Lambda^4 - \phi^4 + \text{h.c.})/2$, then we have $m_a = (f/\Lambda)^n/(n \Lambda^2 f)$. As shown in Fig. 1 the pseudo-Goldstone boson arising from a global symmetry $U(1)_{\text{gl}}$ does not appear in the loops if each vertex satisfies $U(1)_{\text{gl}}$. But, it is known that all global symmetries are approximate \[44,95,98\]. Most strong explicit breaking may be from the anomaly of the type $U(1)_{\text{gl}}G$, where $G$ is a non-Abelian gauge group.

The most wait-for pseudoscalar particle is the very light axion in the axion window because its discovery will confirm at least three: (1) a physical confirmation of instanton solutions of non-Abelian gauge theories \[92\]. (2) ‘t Hooft solution \[100\] of the $U(1)$ problem of QCD \[101\], and (3) at least some portion of CDM in the Universe. The particle axion was first appreciated by Weinberg and Wilczek in the Ben Lee Memorial Conference in October, 1977 \[102\], using the Peccei-Quinn (PQ) symmetry \[103\]. If $G$ is QCD, the symmetry $U(1)_{\text{gl}}$ is called the PQ symmetry $U(1)_{\text{PQ}}$ and the pseudo-Goldstone boson a related to $U(1)_{\text{PQ}}$ is called the QCD axion. The axion is needed to understand the strong CP problem of “Why is the neutron electric dipole moment so small even though the gluon interactions (in the presence of instanton solutions of QCD) allow a neutron-size dipole moment?”. In early days, three kinds of solutions to the strong CP problem were admitted \[11\]: the calculable solution, the massless up quark case, and the axion solution. The calculable solutions have not provided yet an acceptable model with sufficiently small quark electric dipole moment. The massless up quark case is not favored in the global fit \[14\]. The remaining axion solution is checked in various cases as discussed in the next Subsect.

Field theory examples on axions with renormalizable couplings corresponding to BCM1 are usually classified to the KSVZ and DFSZ models \[7,10\]. But, this classification is too simple. There can be many KSVZ and DFSZ type models with one type of quark representations \[104\]. One may introduce many different types of quarks also for axion phenomenology.

Therefore, it is better to have a theory predicting definite PQ charges of the quarks in a full theory. The most attractive proposal along this line is the string compactification. Here, the PQ global symmetry is determined once the compactification scheme is presented. Standard models obtained from string compactification include many quarks beyond the Standard Model spectrum, in particular numerous singlet fields. Along this line, several years ago a QCD axion including non-renormalizable terms was studied and the axion-photon-photon coupling has been calculated with an approximate $U(1)_{\text{PQ}}$ symmetry \[105\]. Recently, an exact $U(1)_{\text{PQ}}$ symmetry has been studied in a string compactification where the axion-photon-photon coupling has been calculated below the PQ symmetry breaking scale \[106\].

$$c_{a\gamma\gamma} = \frac{1123}{388} - 1.98 \approx 0.91.$$ \[7\]

We expect that more calculations of $c_{a\gamma\gamma}$ will be performed in string models with the property of successful SM phenomenologies, which will guide us where to look for the QCD axion \[55\].

Dark energy can be the case of CChmp in the above classification. Pseudoscalar CChmp have been discussed already more than a decade ago in Refs. \[50,54,107\]. But, a more plausible analysis, looking into the detail of string compactification, has been presented recently \[44,98,108\].

The field mass $m_a$ responsible for dark energy corresponds to $m_a \approx 10^{-33}$ eV \[62\]. Meanwhile, if the axion field is responsible for CDM, the typical mass scale is between $10^{-5}$ eV and $10^{-2}$ eV \[13\]. In string theory there are many ultralight axions possibly down to the Hubble scale $H_0 = 10^{-33}$ eV \[109\]. Axions in the mass range between $10^{-28}$ eV and $10^{-18}$ eV become non-relativistic at a later cosmological epoch relative to the standard CDM. Such a light scalar field leads to the suppression of the CDM power spectrum on small scales \[107,111\] (like light massive neutrinos), so there is an observational signature for ultralight axions if the axion potential is of the form \[1 - \cos(a/f_a)]^3 \[108\].

\[2\] For this specific form, one needs fine-tunings between domain wall number one, two, and three terms in the potential.
The expected power conversion \( P_{a\rightarrow\gamma} \) is extraordinarily small, but nonetheless it can be within the present experimental capabilities for an axion mass in the 1–20 \( \mu \)eV range. Have we known the axion mass with a 1 part per million (ppm) accuracy, it would take less than a day to detect it if axions were more than 10% of the DM. The main issue is that, barring the BICEP2 results \[112\]–\[116\], we have no such information. The best-suited axion DM mass is below about 1 meV all the way to about 1 \( \mu \)eV, spanning three orders of magnitude with a potential line width of about 1 ppm. Clearly, scanning the whole axion mass range will require too many steps, and therefore the sensitivity needs to be very high at each step.

Furthermore, in some theoretical scenarios, the axion DM mass is not constrained from below and can be very light, well below 1 \( \mu \)eV. In addition to the microwave cavity method, which is mostly applicable between 1-20 \( \mu \)eV, other methods include looking for axions emitted by Sun’s core, and astrophysical limits, as axions can provide another channel of energy loss, significantly altering the star lifetime. An overview of the present experimental/astrophysical limits of the axion coupling constant vs. the axion mass are given in Fig. 4.

Looking at Eq. (8), it is clear there is a number of possible improvements one can make in this method: (i) Increase the magnetic field value, (ii) Increase the magnetic field volume, and (iii) Increase the cavity quality factor. The pioneering axion DM experiments that started in the late 1980’s \[117\]–\[118\] probed an axion DM candidate in a limited mass region, assuming a stronger axion to photon coupling than is required by theory by roughly two orders of magnitude.

Over a period of more than 15 years, the dominant axion dark matter experiment (ADMX), currently located at the University of Washington and ADMX-HF located at Yale University, have made several conceptual improvements and have improved on those limits. The second generation ADMX experiment, owing to the development of very low noise SQUID amplifiers just below 1 GHz \[119\] and a number of additional smaller developments, has reached the boundaries of a plausible axion DM candidates. Currently implementing a dilution refrigerator to their system is expected to allow them to either detect or exclude an axion comprising 100% of the DM for masses in the range 1–20 \( \mu \)eV.

The new Center for Axion and Precision Physics (CAPP) \[120\], established by the Institute for Basic Science in South Korea \[121\], plans to either detect or exclude an axion DM component down to the 10% level for a similar axion mass range. This will be achieved by \[55\] (a) Development of a 25 T and then a 35 T solenoidal magnet compared to the currently used 8–9 T solenoidal

\[
P_{a\rightarrow\gamma} = g_{a\gamma\gamma}^2 \frac{\rho_a}{m_a} B_0^2 V C_j Q_L ,
\]
magnets, (b) Substantially improving, roughly by an order of magnitude, the quality factor of the microwave cavities in the presence of strong magnetic fields, and (c) Constructing and running a toroidal cavity with a large volume and a reasonable B-field value so that the overall product $B^2V$ is an order of magnitude larger than present values.

The commonly used NbTi superconducting cable has a critical current that falls very rapidly as the magnetic field increases above 10 T, making it unsuitable to obtain higher B-field strengths. However, recent developments with high $T_c$ cables makes possible achieving much higher current densities at large B-field values, when they are cooled at low temperatures around 4 K. This is an experimental method fuelled by the energy-storage field and prototype magnets are already under development. CAPP is collaborating with the Magnet Division of Brookhaven National Laboratory to develop a 10cm inner bore diameter capable of producing around 25 T of magnetic field. Preliminary tests on different high $T_c$ cables are providing encouraging results that the goal can be met. The expected time period for this development is of order of five years, after which we develop a separate magnet with a goal of achieving 35 T peak magnetic field, albeit with smaller inner bore diameter. The next step would be to configure a toroidal magnetic field, optimising the use of the magnetic field as the fringe field is minimized in that geometry. Preliminary cable testing results also point to this geometry for achieving the highest possible magnetic field values. The time scale for this development is of order ten years.

The presently used cavities have a quality factor between 50 K–100 K. It has been reported by ADMX that they are developing cavities with thin-film superconducting coatings on the vertical side walls with the goal of increasing the cavity quality factor by roughly a factor of five. This is possible when the B-field is shaped to be aligned with the vertical wall, minimising the transverse B-field below about 100 Gauss. Further increases of the quality factor are hindered by the top/bottom surfaces of the right-cylindrical cavity as the magnetic field angle traversing the surface is very close to 90 degrees. Our plan to further improve upon this achievement is two-fold: First, develop a toroidal cavity where the B-field can be shaped along the cavity walls reducing the transverse B-field below the required level. If that is possible, the quality factor can be increased by several orders of magnitude. Second, the top/bottom plates are going to be treated in a way that the B-field can penetrate it without affecting the superconducting layer on the inside of the cavity. Again, the quality factor can potentially increase by several orders of magnitude.

The expected axion width is of order 1 ppm, i.e. the axion quality factor is a bit better than $10^6$. Therefore, the best one can do is to produce a cavity with the same quality factor, so the best one can expect is to gain a factor of 10 to 20 in the axion to photon power conversion. The scanning speed goes as the square root of the quality factor since there are more steps required in order to cover all possible frequencies, i.e. the best one can expect to do is a scanning speed improvement factor between three and five.

BICEP2 results favour axion masses in the meV range, albeit with only 1–10% of DM composed of axions. This fact makes it particularly difficult to detect it as the volume of microwave cavities are particularly small and not of much practical use at those frequencies, plus the axion DM density is very weak. If the BICEP2 results turn out to be confirmed,$^3$ one could follow a different strategy in detecting axions$^{122}$. If the axion mass were to be found, then one could launch a dedicated axion DM experiment within a very small axion mass range having much higher chances of success.

$^3$ We note that the recent Planck-dust report extrapolated to the BICEP2 field gives the dust contribution similar to $r \approx 2$ without the dust contribution$^{123}$.
III. THE COSMOLOGICAL CONSTANT PROBLEM AND STRING THEORY

In order to realize the present-day cosmic acceleration with the cosmological constant $\Lambda$, we require that $\Lambda$ is of the order of $H_0^2$, i.e., $\Lambda \approx H_0^2 = (2.1332h \times 10^{-45} \text{ GeV})^2$, where $h \approx 0.7$. This corresponds to the energy density $\rho_\Lambda \approx \Lambda M_p^4 \approx 10^{-120} M_p$. Even before the discovery of the present-day cosmic acceleration, Weinberg [24] put the bound on $\rho_\Lambda$, as

$$-2 \times 10^{-120} M_p^4 \lesssim \rho_\Lambda \lesssim 6 \times 10^{-118} M_p^4. \quad (9)$$

The lower bound comes from the fact that the negative cosmological constant does not lead to the collapse of the Universe today. The upper bound corresponds to the requirement that the vacuum energy does not dominate over the matter density for redshifts $z$ larger than 1 to realize the successful structure formation.

There have been attempts to explain the very low values of $\rho_\Lambda$ ranging the Weinberg bound (9). For example, Bousso and Polchinski [12] introduced the 4-form field $F_{\mu\nu\lambda\sigma}$ with the energy density $F_{\mu\nu\lambda\sigma} F^{\mu\nu\lambda\sigma}/48 = c^2/2$, where $c$ is a constant. In the context of string theory, there are “electric charges” (membranes) sourcing the 4-form field dual to “magnetic charges” (5-branes). The constant $c$ can be quantized in integer ($n$) multiples of the membrane charge $q$, such that $c = nq$.

Bousso and Polchinski introduced $J$ 4-form fields together with $J$ membrane species with charges $q_1, q_2, \cdots, q_J$. The number $J$ can be as large as 100 in string theory. Since the flux energy density of each charge is given by $n_i^2 q_i^2/2$, the effective cosmological constant reads

$$\Lambda = \Lambda_b + \sum_{i=1}^J n_i^2 q_i^2/2, \quad (10)$$

where $\Lambda_b$ is the bare cosmological constant. For the anti de Sitter minimum with $\Lambda_b < 0$, there exist integers $n_i$ satisfying

$$2|\Lambda_b| \lesssim \sum_{i=1}^J n_i^2 q_i^2 < 2(|\Lambda_b| + \Delta \Lambda), \quad (11)$$

where $\Delta \Lambda \simeq 10^{-123}$ in the unit $M_p = 1$.

If we consider a $J$-dimensional grid with axes corresponding to $n_i q_i$, the displacement of the 4-field form is given by discrete grid points with integers $n_i$. The region $i$ corresponds to a thin-shell characterized by the radius $r = \sqrt{2|\Lambda_b|}$ and the width $\Delta r = \Delta \Lambda/\sqrt{2|\Lambda_b|}$. The volume of the thin-shell is

$$V_S = \Omega_{J-1} r^{J-1} \Delta r = \Omega_{J-1} |2\Lambda_b|^{J/2-1} \Delta \Lambda, \quad (12)$$

where $\Omega_{J-1} = 2\pi^{J/2}/\Gamma(J/2)$ is the area of a unit ($J-1$)-dimensional sphere. A grid cell has a volume $V_C = \prod_{i=1}^J q_i$. There is at least one value of $\Lambda$ for $V_C < V_S$, i.e.,

$$\prod_{i=1}^J q_i < \frac{2\pi^{J/2}/\Gamma(J/2)}{2\Lambda_b}^{J/2-1} \Delta \Lambda. \quad (13)$$

When $J = 100$, $|\Lambda_b| = 1$ and $\Delta V = 10^{-123}$ with equal charges ($q_i = q$, for $i = 1, 2, \cdots, J$), the condition (13) is satisfied for $q < 0.035$. Since the charge $q\sqrt{\Lambda}$ has the dimension of mass from Eq. (10), this condition translates to $\sqrt{\Lambda} < 0.19$ in units of $M_p$. Thus, the presence of many 4-form fields allows the possibility of realizing a small effective cosmological constant.

The idea of Bousso and Polchinski is based on the flux energy density originating from multiple 4-form fields. This idea was extended to the so-called flux compactification on a Calabi-Yau manifold in type II string theory. In the presence of fluxes, Kachru, Kallosh, Linde and Trivedi [12] first set up a supersymmetric anti de Sitter (AdS) vacuum with all moduli fields fixed. Then, they obtained a de Sitter vacuum by adding an anti D3-brane in a warped geometry to lift up the AdS state.

There are hundreds of different 3-cycles on the Calabi-Yau manifold in the flux compactification. A macroscopic observer can view a 5-brane wrapping a 3-cycle as a 2-brane (membrane). The 5-brane can wrap any of these 3-cycles, which gives rise to hundreds of different membranes in four-dimensional space-time. The number of vacua appearing in string theory can be extremely large. For 500 three-cycles with each cycle wrapped by up to 10 fluxes, we have $10^{500}$ vacua.

The possible presence of such a large amount of vacua led to the notion of so-called string landscape [126]. This landscape includes so many possible configurations of local minima, among which our Universe may correspond to one of them. Each vacuum in the string landscape has different matter and coupling constant. The SM is not predicted uniquely in this picture. The argument is that we may be able to find a vacuum with an extremely small energy density among $10^{500}$ vacua. However, this anthropic argument depends on “Those packed near $\Lambda = 0$ out of $10^{500}$ vacua describe particle phenomenology correctly, in particular with three chiral families and $\sin^2 \theta_W = 3/8,”$ otherwise the landscape vacua differing by $\Delta \Lambda$ describe unacceptable universes. From this reasoning, the string landscape is commented in Table I as ‘not yet’ established.

A general problem with the anthropic arguments is that they are often applied to a single parameter while fixing all the others. A parameter value that is ruled out in one case may be acceptable if something else is changed at the same time. In this sense, it is not clear that the anthropic arguments of $\Lambda$ provide a satisfactory answer to the cosmological constant problem.

As commented before, the DE scale may be accountable from highly suppressed non-renormalizable terms in string-allowed discrete symmetries [44, 98] if the true vacuum has zero cosmological constant. In this sense,
the theoretical solution toward the vanishing cosmological constant is more difficult to solve than obtaining a tiny DE scale on top of the vanishing cosmological constant $23, 26$.

IV. INFLATION

The possibility of an exponential expansion of the Universe was known even before the influential paper of Guth [45] which advocates diluting away the GUT scale monopoles [132]. For example, in the abstract of the Kazanas’s paper [130], it is stated that “...In particular it is shown that under certain conditions this expansion law is exponential. It is further argued that under reasonable assumptions for the mass of the associated Higgs boson this expansion stage could last long enough to potentially account for the observed isotropy of the universe.” In the papers of Sato [128, 129], diluting away topological defects such as monopoles and domain walls was stressed after the advent of the modern GUT model [133, 134]. In the Guth’s paper [45] it was clearly emphasized that “...In particular it is shown that under certain conditions this expansion law is exponential. It is further argued that under reasonable assumptions for the mass of the associated Higgs boson this expansion stage could last long enough to potentially account for the observed isotropy of the universe.” In the papers of Sato [128, 129], diluting away topological defects such as monopoles and domain walls was stressed after the advent of the modern GUT model [133, 134].

The scalar field responsible for inflation is called ‘inflaton’. The inflaton field is a superposition of quanta of all possible wave lengths. A quantum fluctuating scale inflates exponentially and after passing the horizon, it is stretched exponentially with an almost scale-invariant form [135–139] and the frozen-scale still inflates exponentially (see Ref. [140] for a review). Different fluctuating scales go out of the horizon at different cosmic times and their exponentially stretched scales are correlated.

After the end of inflation, the quantum fluctuations enter the horizon again and become the sources of density perturbations. The prediction of nearly scale-invariant primordial perturbations generated during inflation was consistent with the temperature anisotropies of Cosmic Microwave Background (CMB) observed by the COBE satellite [141]. The recent WMAP and the Planck data of CMB refined the temperature anisotropies to very high accuracy [4, 12].

The observables and the constraints implied by inflation are

- A sufficient inflation, requiring the large e-fold number, $N_e > 70$, for addressing horizon and flatness problems.
- The amplitude of temperature anisotropies $δT/T \approx 10^{-5}$, for galaxy formation with CDM.
- The spectral index of scalar perturbations $n_s \approx 0.96$, from WMAP and Planck data.

As long as the slow-roll conditions are satisfied, the single-field inflationary scenario generally gives rise to local non-Gaussianities with $f_{\text{NL}}$ much smaller than 1 even for most general scalar-tensor theories with second order equations of motion [143–145]. Hence the slow-roll single-field models are consistent with the Planck bound of non-Gaussianities. Using the observational bounds of $n_s$ and $r$, we can distinguish between many single-field inflationary models [146–148]. For example, the self-coupling potential $V(\phi) = \lambda \phi^4/4$ [149] and hybrid inflation [150] with $n_s > 1$ are disfavored from the data.

The amplitude of tensor perturbations is given by $P_T = 2H^2/(\pi^2 M_P^2)$, so the detection of gravitational waves in CMB observations implies that the energy scale of inflation is directly known [151–153]. Since the B-mode polarization of CMB is generated by tensor perturbations but not by scalar perturbations, the B-mode detection is a smoking gun for the existence of primordial gravitational waves.

If the tensor-to-scalar ratio $r$ is smaller than the order of 0.01, it is not easy to detect the CMB B-mode polarization. If $r$ is detected in the range $r > 0.05$, then the energy scale during inflation corresponds to the GUT scale. The great interest in the announcement of $r \sim 0.16$ from the BICEP2 group [62] is because of the implication that the Universe once passed the vacuum energy scale of $10^{16}$ GeV. Even though the GUT scale $M_{\text{GUT}}$ is humongous from our TeV scale Standard Model, it is tiny from the point of gravity scale, the Planck mass $M_P$.

- The tensor-to-scalar ratio $r \lesssim 0.2$, from WMAP and Planck data.
- The non-linear estimator of scalar non-Gaussianities for the local shape $f_{\text{NL}} = 2.7 \pm 5.8 (68\% \text{ CL})$, from Planck data.

4 From the Planck data the existence of CDM was also confirmed (by 7σ [142]) better than any other data.

5 However, the future observations like LiteBIRD may reach this range.
Because of the micro density perturbation, the vacuum energy at the scale \(10^{16} \text{ GeV}^4\) leads to \(r \sim O(0.1)\). This phenomenon of the GUT scale energy density during inflation is usually parametrized by chaotic inflation with the potential \(V(\phi) = \frac{1}{4} m^2 \phi^2\) [149].

If a large \(r\) of order 0.2 is detected, the field value in the \(\phi^2\) chaotic inflation is bounded from below, i.e. \(\langle \phi \rangle > 15M_P\), which is known as the ‘Lyth bound’ [150]. This situation is shown in Fig. 6 where the energy density at the inflationary epoch is the GUT scale. The field value \(\langle \phi \rangle > 15M_P\) is trans-Planckian and the GUT energy density at \(M_P\) is tiny. So, one needs a fine-tuning in the \(\phi^2\) chaotic inflation. Introducing a confining force at a GUT scale, a heavy axion for the inflaton with a potential bounded from above was proposed [157], which is called natural inflation. In this scenario, the energy density has the upper bound of order \(M^4_{\text{GUT}}\) as shown in Fig. 6(a). One period of the inflaton in this case is of order \(M_P\), and hence the Lyth bound is violated. To remedy this, two confining forces are introduced with two heavy axions with the resulting potential [158],

\[
V = -\Lambda^2_1 \cos \left( \frac{\alpha_1}{F_1} + \frac{\beta_2}{F_2} \right) - \Lambda^2_2 \cos \left( \frac{\alpha_1}{F_1} + \frac{\delta_2}{F_2} \right) + \text{constant},
\]

where \(\alpha, \beta, \gamma, \text{and} \delta\) are determined by the corresponding PQ symmetries of two heavy axions \(a_1\) and \(a_2\). Even though we allow \(O(1)\) couplings, the GUT mass scales can lead to \(M_P\) with the probability of \(~1\%\). With mass parameters of \(50M_{\text{GUT}}\), we would obtain \(50M_P\) with the probability of \(~1\%\). This is the Kim-Nilles-Peloso 2-flation model [158]. The probability of the 2-flation model with a large decay constant, i.e. \(f_\phi > 15M_P\) to occur as shown in Fig. 6(b), is about \(1\%\). The green-potential in Fig. 6(b) is the other heavy axion potential. It can be generalized to N-flation [159].

The axionic topological defects in the anthropic window [25, 160] can be diluted away if inflation occurs below the anthropic window scale. With the GUT scale energy density during inflation, however, this dilution mechanism does not work. With the GUT energy scale inflation as implied by the BICEP2 [62], it could have pinned down to \(f_\phi \sim 10^{11} \text{ GeV}\) [115, 116], using the numerical calculation of radiating axions from axionic string-wall system [161]. In the numerical calculation, the Vilenkin-Everett mechanism [162] of erasing the horizon scale string has not been taken into account. In addition, the hidden-sector confining force can erase horizon scale axionic strings such that the QCD axion domain wall is not a serious cosmological problem [163]. The hidden-sector solution needs the hidden-sector domain-wall number of \(N_0 = 1\), which is possible in string compactification with an anomalous \(U(1)\) [164].

In addition to pinning down the upper bound on \(f_\phi\), the GUT scale inflation provokes a question, “What is the symmetry which naturally satisfies the Lyth bound?” Lyth considered this problem with respect to the \(\eta\) parameter [165]. But, there exists a more fundamental question. In an ultra-violet completed theory such as string theory, every parameter is calculable. If we consider the \(\phi^2\) chaotic inflation of Fig. 6 there is a question, “Why do we neglect other terms?” In string theory, only discrete symmetries are permitted by the compactification process. For example, a term \(\phi^{103}/M_P^{100}\) can be possible if the discrete symmetries allow it. But with the trans-Planckian value, for example \(\langle \phi \rangle \sim 31\), the coefficient must be tuned to 1 out of \(10^{127}\), which is as bad as the cosmological constant problem.

Fortunately, there is another way for inflation to occur. We must choose the hilltop inflation, but sacrificing the single-field inflaton. It is not so bad in view of the fact that the 2-flation model already introduced two axions in the inflaton sector. Then, the inflating region is near origin such that the minimum at \(f_{DE}\) is far away from the origin. In the region \([0, f_{DE}]\) the vacuum energy is of order \(M^4_{\text{GUT}}\). This can be obtained from the condition on the quantum numbers of the assumed discrete symmetry [166]. The inflaton rolls in the yellow region in Fig. 7 where the inflaton takes a green curve in the two-inflatons space.
V. DISCUSSION

After the discovery of the fundamental Brout-Englert-Higgs boson, which is represented as $H_u$ and $H_d (= H^0)$ in non-SUSY case, we reviewed the cosmological role of spin-0 bosons. This finding hints the possibility of numerous spin-0 bosons ($\phi$) at the GUT scale. Spin-0 bosons at the GUT scale of the canonical dimension 1 can have more important effects to low energy physics compared to those of spin-$\frac{1}{2}$ fermions of the canonical dimension $\frac{1}{2}$ (Dirac fermions $\psi, \bar{\psi}$ for example) at the GUT scale. For example, the spin-0 contribution

$$\frac{\phi^{2n}}{M_p^{2n+2m}}(H_d H_u)^m$$

(15)

dominates the fermion contribution $\frac{\bar{\psi}^n \psi^n}{M_p^{n+2m}}(H_d H_u)^m$ for $n, m \geq 1$. In addition, the existence of fundamental spin-0 bosons at the GUT scale may be extended to a larger symmetry: supersymmetric GUTs, or minimal supersymmetric Standard Models from string compactification. The interactions of the singlet fields only can take a SUSY superpotential, for example with GUT scale singlets $\phi$ and trans-Planckian singlets $\Phi$ for simplicity [16],

$$W = \sum_i \frac{\phi^{a_i}}{M_p^{a_i+\ell_i-3}}\Phi^{\ell_i}. \quad (16)$$

The rationale leading to the forms of Eqs. (15) and (16) are the discrete symmetries obtained from string compactification [44] which guarantees the absence of gravity obstruction of discrete symmetries, for example via wormholes [95]. The form of the interaction (16) can lead to inflation with trans-Planckian decay constant with a multi-field hilltop potential, i.e. BCM2. The form of the interaction (15) can lead to QCD axion via BCM1, and the DE scale via CCM. The fundamental scalars at the TeV, GUT and trans-Planckian scales allow all scenarios presented in Subsec. [11][A] These are worked out on top of vanishing cosmological constant, which is assumed in any particle physics models. At present, we do not have any persuasive hint toward a theoretical solution of the vanishing cosmological constant. Any theory for the vanishing cosmological constant must satisfy the requirements of particle phenomenology we used in this review.

The fundamental scalars may be detectable if their couplings to gluons are appreciable. The front runner in the search of fundamental scalars hinting high energy (GUT or intermediate) scales is the QCD axion which couples to the gluon anomaly.

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