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Axiogenesis

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We propose a mechanism called axiogenesis where the cosmological excess of baryons over antibaryons is generated from the rotation of the QCD axion. The Peccei-Quinn (PQ) symmetry may be explicitly broken in the early Universe, inducing the rotation of a PQ charged scalar field. The rotation corresponds to the asymmetry of the PQ charge, which is converted into the baryon asymmetry via QCD and electroweak sphaleron transitions. In the concrete model we explore, interesting phenomenology arises due to the prediction of a small decay constant and the connections with new physics at the LHC and future colliders and with axion dark matter.

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Introduction.—One of the goals of fundamental physics is to understand the origin of the Universe. For this purpose, the standard model (SM) of particle physics needs an extension to explain the cosmological excess of matter over antimatter. Mechanisms to generate the baryon asymmetry have been intensively studied in the literature under the name of baryogenesis. The proposed origins of the baryon asymmetry include explicit baryon or lepton number violation from (i) the supersymmetric partners of baryons or leptons in the Affleck-Dine mechanism [1,2], (ii) anomalous baryon number violating processes in electroweak baryogenesis [2–6], and (iii) heavy right-handed Majorana neutrinos in leptogenesis [7,8]. Developing novel baryogenesis mechanisms has been one of the main focuses of particle physics in the past decades.

The SM also needs an extension to explain the smallness of charge conjugation parity (CP) violation in QCD [9] which on theoretical grounds is expected to be large [10]. This is known as the strong *CP* problem and can be elegantly solved by the Peccei-Quinn (PQ) mechanism [11,12]. The so-called PQ symmetry is spontaneously broken to yield a pseudo Nambu-Goldstone boson, the axion [13,14]. The PQ symmetry is explicitly broken by the quantum effects of QCD of the Adler-Bell-Jackiw type [15,16]. The quantum effects give a potential to the axion and drive the axion field value to the point where *CP* symmetry is restored, solving the strong *CP* problem. The PQ mechanism is especially attractive because the axion is also a dark matter candidate [17–19], which provides yet another missing piece of the standard model.

We discover that when the PQ mechanism is introduced into the SM, the baryon (B) and lepton (L) asymmetries are generated in a wide class of models. We call the following baryogenesis scheme as axiogenesis, which in general includes two main ingredients: (i) an asymmetry of the PQ charge is generated in the early Universe as a coherent rotation in the axion direction and (ii) the PQ asymmetry is later transferred to the $B+L$ asymmetry via the QCD and electroweak sphaleron transitions. (We may convert the $B+L$ asymmetry into the $B-L$ asymmetry by some $B-L$ breaking interaction. Such a scenario will be investigated in a future work [20].) We contrast axiogenesis with other existing baryogenesis models after we introduce a concrete example.

The PQ symmetry is an approximate global symmetry which is explicitly broken by the QCD anomaly. Given that the symmetry is not exact, it is plausible that the PQ symmetry is significantly broken in the early Universe, and the rotation of the axion is induced. In fact, it is expected that quantum gravity does not allow for a global symmetry [21–25] and the PQ symmetry is at best understood as an accidental symmetry explicitly broken by higher dimensional operators [26–29]. Even when one requires that this explicit breaking not spoil the solution to the strong *CP* problem in the present Universe, the rotation can still be induced from such interactions in the early Universe as we will describe. Another example is a larger QCD scale in the early Universe [30–33], which can initiate the axion oscillation and, once the QCD scale becomes small enough, the axion begins to rotate. These PQ-breaking sources well justify the axion rotation.

A fast rotation of the axion corresponds to a large PQ charge asymmetry. The PQ symmetry and the SM quark chiral symmetries are explicitly broken by the quantum effect of QCD, called the QCD anomaly. In the thermal bath of the early Universe, a nonperturbative process called the QCD sphaleron transition is active. The transition, through the anomaly, converts the PQ charge asymmetry into the quark chiral asymmetry until the asymmetries reach equilibrium values. The quark chiral symmetry and the $B+L$ symmetry are also explicitly broken by a weak anomaly. The quark chiral asymmetry is then converted into the $B+L$ asymmetry by another nonperturbative process known as the electroweak sphaleron transition. We may also consider a model with a weak anomaly of the PQ symmetry, as is the case with the KSVZ model [34,35] embedded into grand unification and the supersymmetric DFSZ model [36,37]. In such a model the PQ asymmetry is directly converted into the $B+L$ asymmetry via electroweak sphaleron transitions. Consequently, the rotation of the axion can account for the observed matter asymmetry of the Universe via the QCD and electroweak sphaleron transitions.

Baryon asymmetry from axion rotation.—We discuss a minimal version of axiogenesis that achieves the conversion between the PQ asymmetry $n_{\rm PQ}$ in the form of the axion rotation $\dot{\theta}$ and the baryon asymmetry solely by the SM QCD and electroweak sphaleron processes.

The axion ϕ_a is the angular direction of the complex scalar field

$$P = \frac{1}{\sqrt{2}}(S + f_a)e^{i(\phi_a/f_a)}, \tag{1}$$

whose radial direction obtains a vacuum expectation value f_a, which is called the axion decay constant, and breaks the PQ symmetry. Analogous to how classical rotational symmetry leads to angular momentum conservation, the shift symmetry $\phi_a \to \phi_a + \alpha f_a$ implies a conserved Noether charge associated with the rotation in the axion direction. The PQ charge asymmetry $n_{\rm PQ}$ is exactly the Noether charge density associated with the shift symmetry. We define $n_{\rm PQ}$ with the following normalization, $n_{\rm PQ} = iP\dot{P}^* - iP^*\dot{P}$, where the dot denotes a time derivative. When the radial mode is settled to the minimum f_a, the PQ charge asymmetry is then given by

$$n_{\rm PQ} = \dot{\theta} f_a^2, \tag{2}$$

where $\theta \equiv \phi_a/f_a$. Here we simply assume the rotation exists, while we present a concrete model to initiate the axion rotation in the next section.

The PQ asymmetry is converted into chiral asymmetries of SM quarks via QCD sphaleron transitions. The chiral asymmetries are then converted into the $B+L$ asymmetry via electroweak sphaleron transitions. Although the chiral symmetries are explicitly broken by the SM Yukawa couplings and hence the asymmetries are constantly washed out, the large PQ asymmetry continuously sources the chiral asymmetries and a nonzero baryon asymmetry remains in a quasiequilibrium state. If the PQ symmetry has a weak anomaly, the PQ asymmetry is directly converted into $B+L$ asymmetry. In short, the PQ asymmetry is converted into $B+L$ asymmetry by QCD and electroweak sphaleron transitions. With the detail given in the Supplemental Material [38], we find that, before the electroweak phase transition, the baryon number density n_B is given by

$$n_B = c_B \dot{\theta} T^2, \qquad c_B \simeq 0.1 - 0.15 c_W. \tag{3}$$

Here c_W is the weak anomaly coefficient of the PQ symmetry normalized to that of the QCD anomaly. The electroweak sphaleron process becomes ineffective after the electroweak phase transition and the baryon asymmetry is frozen. The resultant asymmetry normalized by the entropy density s is

$$Y_B = \frac{n_B}{s} = \frac{45 c_B}{2 g_* \pi^2} \frac{\dot{\theta}}{T}\bigg|_{T=T_{\rm ws}} \simeq 2\times 10^{-3}\left(\frac{c_B}{0.1}\right)\frac{\dot{\theta}(T_{\rm ws})}{T_{\rm ws}}, \tag{4}$$

where $T_{\rm ws}$ is the temperature below which the electroweak sphaleron transition becomes ineffective and g_* is the effective degrees of freedom in the thermal bath.

For $\dot{\theta}$ required to reproduce the baryon asymmetry, the axion continues to rapidly rotate even around the QCD phase transition. Even when the axion mass becomes comparable to the Hubble expansion rate, the oscillation does not occur because the kinetic energy of the rotation is still much larger than the barrier of the axion cosine potential. The actual oscillation around the minimum is delayed until when the kinetic energy becomes comparable to the potential energy of the axion. Therefore, the axion abundance becomes enhanced [39] in comparison with the conventional misalignment mechanism [17–19].

As derived in the Supplemental Material [38], assuming PQ charge conservation, $\dot{\theta}$ is a constant before the PQ breaking field reaches the minimum, whereas $\dot{\theta} \propto a^{-3}$ thereafter, with a the scale factor. Assuming the latter case at the weak scale, we find the axion abundance

$$\frac{\Omega_a h^2}{\Omega_{\rm DM} h^2} \simeq 140\left(\frac{f_a}{10^8\ {\rm GeV}}\right)\left(\frac{130\ {\rm GeV}}{T_{\rm ws}}\right)^2\left(\frac{0.1}{c_B}\right), \tag{5}$$

to be much larger than the observed DM abundance $\Omega_{\rm DM}h^2$ for f_a satisfying the astrophysical constraints [40–46], the SM prediction $T_{\rm ws} \simeq 130$ GeV [47], and $c_B = \mathcal{O}(0.1\text{–}1)$. (A value of $f_a = \mathcal{O}(10^6)$ GeV leads to both successful axiogenesis and axion dark matter and interestingly resides in the so-called axion hadronic window [42,48,49], which however is recently under scrutiny [50,51].) We require either (i) the axion rotation is damped after the electroweak phase transition, (ii) the electroweak phase transition occurs earlier than the SM prediction, or (iii) $c_B \gg \mathcal{O}(1)$ because of a large coefficient of the weak anomaly.

When the Higgs couples to particles with masses above the electroweak scale, it is possible that the electroweak phase transition occurs at a high temperature, and the Higgs eventually relaxes to the electroweak scale. We present a toy model in the Supplemental Material [38].

A large weak anomaly coefficient is possible in multi-field extensions of the Kim-Nilles-Peloso mechanism [52–57], as considered in [58]. Assuming axion dark matter, the axion-photon coupling is

$$|g_{a\gamma\gamma}| = \frac{\alpha(c_W + c_Y)}{2\pi f_a} \simeq 10^{-9}\ \mathrm{GeV}^{-1}\left(\frac{130\ \mathrm{GeV}}{T_{\mathrm{ws}}}\right)^2 \quad (6)$$

where α is the fine structure constant. This prediction assumes that the hypercharge anomaly coefficient c_Y of the PQ symmetry is negligible. For $T_{\mathrm{ws}} = 130$ GeV, this large coupling is excluded by the limit from CAST [59], $|g_{a\gamma\gamma}| < 6.6 \times 10^{-11}\ \mathrm{GeV}^{-1}$. However, the contribution from the hypercharge anomaly can reduce or even exactly cancel the coupling.

We treat the rotation as a background field. A small portion of the PQ asymmetry is converted into the quark chiral asymmetries which are washed out by the Yukawa couplings. The washout interaction is suppressed by a small up quark Yukawa coupling y_u because in the limit of a vanishing y_u, a linear combination of the PQ symmetry and the up quark chiral symmetry is exact and washout does not occur. As is shown in the Supplemental Material [38], the washout of the PQ asymmetry is negligible.

We comment on the similarities and the differences of axiogenesis with the models in the literature. In spontaneous baryogenesis [60,61], baryon asymmetry is generated by a chemical potential of baryons given by the motion of a pseudo Nambu-Goldstone boson. The chemical potential is provided by the oscillation or the slow motion of the boson field driven by an explicit symmetry breaking potential. In axiogenesis, explicit breaking is effective only at higher energy scales and drives the rapid rotation of the axion instead. As a result, axiogenesis is compatible with the QCD axion. Also, in spontaneous baryogenesis the oscillation itself washes out the PQ asymmetry, and the $B+L$ asymmetry needs to be converted into $B-L$ asymmetry, e.g., by the seesaw operator, which is not required in axiogenesis. Baryogenesis using the chemical potential provided by the rotation of the QCD axion is mentioned in [62] but the conversion of the PQ asymmetry into the $B+L$ asymmetry by the QCD and/or weak anomaly is not considered. Baryogenesis via the oscillation of the (QCD) axion by a large mass, the weak anomaly of the PQ symmetry and the seesaw operator is considered in [63]. Baryogenesis by the chemical potential of the weak Chern-Simon number is utilized in the local electroweak baryogenesis [64,65] and other models in [66–68], where the chemical potentials are provided by the Higgs fields and the gluon condensation, respectively.

Affleck-Dine axiogenesis.—In this section we continue the investigation of a concrete realization of axiogenesis by evaluating $\dot{\theta}$. To generate the PQ asymmetry, we employ the idea of Affleck-Dine [1] proposed in a supersymmetric theory, even though supersymmetry is not essential to axiogenesis. (See [69] for a non-supersymmetric Affleck-Dine mechanism.) For clarity and simplicity, we demonstrate a working example by the quartic potential

$$V = \lambda^2\left(|P|^2 - \frac{f_a^2}{2}\right)^2, \qquad \lambda^2 = \frac{1}{2}\frac{m_S^2}{f_a^2}, \quad (7)$$

where P is the complex field breaking the PQ symmetry in the vacuum and m_S corresponds to be the vacuum mass of the radial mode S, which we call the saxion although we do not assume supersymmetry here. The angular mode in the vacuum is the axion. We assume a large initial field value $|P_i| = S_i/\sqrt{2} \gg f_a$, which can arise for a sufficiently small quartic coupling, namely a flat potential of S. (A flat potential is natural in supersymmetric theories, with which we demonstrate axiogenesis using a concrete model and cosmological evolution in the Supplemental Material [38].) The potential at large field values is dominated by the quartic term and thus the saxion mass is initially given by $\sqrt{3}\lambda S_i$. The saxion starts oscillating when the Hubble friction drops below the mass, $3H \simeq \sqrt{3}\lambda S_i$, at the temperature

$$T_{\mathrm{osc}} = \left(\frac{30}{\pi^2 g_*}\right)^{1/4}\sqrt{\lambda M_{\mathrm{Pl}} S_i}, \quad (8)$$

with $M_{\mathrm{Pl}} = 2.4 \times 10^{18}$ GeV the reduced Planck constant.

For large S_i, a higher dimensional potential term that explicitly breaks the PQ symmetry,

$$V_{\not{PQ}} = \frac{P^n}{M^{n-4}} + \mathrm{H.c.}, \quad (9)$$

can be effective. Here M is a dimensionful constant. The potential drives P in the angular direction and causes a rotation. After initiating the rotation, as S decreases by redshifting, explicit PQ breaking quickly becomes very suppressed as it originates from a higher dimensional operator. As a result, the PQ charge becomes conserved soon after the initial motion. It is convenient to normalize the asymmetry by the number density of the saxion,

$$\frac{n_{\mathrm{PQ}}}{n_S} \equiv \epsilon, \quad (10)$$

because this is a redshift-invariant quantity. The scaling of $n_{\mathrm{PQ}} \propto a^{-3}$ can be understood as a result of PQ charge conservation. We use ϵ to parametrize the amount of PQ breaking that leads to the axion rotation or equivalently the potential gradient in the angular direction relative to that of the radial mode at $S = S_i$. We treat ϵ as a free parameter in what follows. In supersymmetric theories described in the Supplemental Material [38], ϵ is naturally order unity.

The saxion acquires a large energy density due to its initial condition $S_i \gg f_a$. While the saxion condensate will eventually thermalize with the SM plasma, as is shown in the Supplemental Material [38], the PQ charge asymmetry is conserved up to cosmic expansion. In other words, thermalization only depletes the energy density in the radial mode and preserves that in the angular mode. Therefore, the rotation continues even after thermalization. Whether the saxion condensate dominates the energy density of the Universe before being thermalized into the SM plasma leads to two possibilities for the subsequent cosmology, both of which we investigate in order.

If the Universe stays radiation-dominated throughout the evolution, the PQ asymmetry due to the axion rotation in units of the entropy density is a redshift-invariant quantity after the onset of the oscillation and is given by

$$Y_{\rm PQ} = \frac{\epsilon V(P_i)}{s m_S(P_i)} = \epsilon \left(\frac{\pi^2 g_*}{30}\right)^{3/4} \frac{15\sqrt{3}}{8\pi^2 g_*} \frac{S_i^{3/2}}{\sqrt{\lambda} M_{\rm Pl}^{3/2}}, \tag{11}$$

which, with Eq. (2), implies that the angular speed is

$$\dot{\theta}(T) = \epsilon \left(\frac{\pi^2 g_*}{30}\right)^{3/4} \frac{S_i^{3/2} T^3}{4\sqrt{3}\sqrt{\lambda} M_{\rm Pl}^{3/2} f_{\rm eff}^2(T)}, \tag{12}$$

where $f_{\rm eff}(T)$ is the effective axion decay constant at temperature T, i.e., $\sqrt{2}|P(T)|$. Finally, based on Eq. (4), the baryon asymmetry is evaluated at the temperature when the electroweak sphaleron is out of equilibrium and reads

$$\begin{aligned} Y_B &= \epsilon \left(\frac{\pi^2 g_*}{30}\right)^{3/4} \frac{15\sqrt{3}}{8\pi^2 g_*} \frac{c_B S_i^{3/2} T_{\rm ws}^2}{\sqrt{\lambda} M_{\rm Pl}^{3/2} f_{\rm eff}^2(T_{\rm ws})} \\ &\simeq 9 \times 10^{-11} \epsilon \xi \left(\frac{S_i}{M_{\rm Pl}}\right)^{3/2} \left(\frac{10^9\ {\rm GeV}}{f_{\rm eff}(T_{\rm ws})}\right)^{3/2} \left(\frac{\rm TeV}{m_S}\right)^{1/2}, \end{aligned} \tag{13}$$

where

$$\xi \equiv \left(\frac{c_B}{100}\right) \left(\frac{T_{\rm ws}}{130\ {\rm GeV}}\right)^2. \tag{14}$$

As illustrated in the previous section, the dark matter abundance in Eq. (5) demands a transfer of the PQ to baryon asymmetries more efficient than that from the SM prediction of $T_{\rm ws} = 130$ GeV along with $c_B \simeq \mathcal{O}(0.1\text{–}1)$ from a weak anomaly coefficient c_W of order unity. This is manifest in the parameter ξ, and theories with a large c_B and/or $T_{\rm ws}$ are also discussed previously.

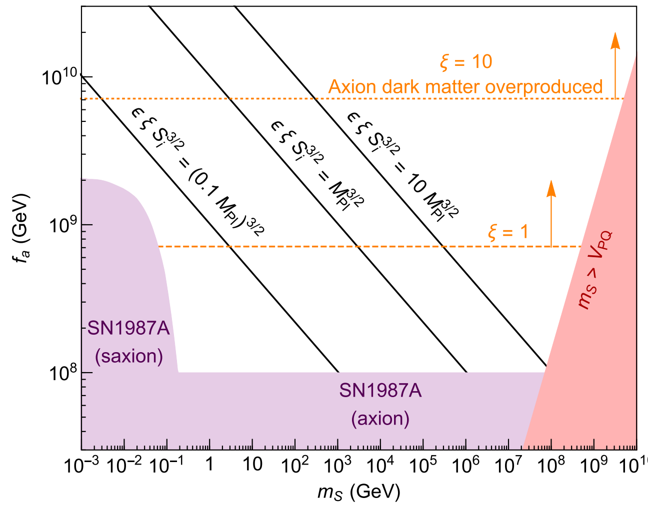

FIG. 1. The parameter space compatible with the observed baryon asymmetry.

In Fig. 1, we demonstrate the viable parameter space of the saxion mass m_S and the axion decay constant f_a. The black contours show the values of $\epsilon\xi S_i^{3/2}$ required by Eq. (13) for the observed baryon asymmetry $Y_B^{\rm obs} = 8.7 \times 10^{-11}$ [70]. From Eq. (5), the region above the orange line is excluded due to axion dark matter overproduction for $\xi = 1$ (dashed) and $\xi = 10$ (dotted). In the red region, the saxion mass m_S exceeds the unitarity limit. The purple region is excluded since the emission of saxions or axions in a supernova core affects the duration of the neutrino emission [40–44,50,51,71]. The constraint from the saxion emission can be, however, evaded by introducing a large enough saxion-Higgs mixing to trap saxions inside the core.

If the saxion dominates, since the P oscillation until thermalization at temperature $T_{\rm th}$, the PQ charge number density $n_{\rm PQ}$ and the saxion number density n_S redshift the same way. After the saxion is depleted to create a thermal bath with a temperature $T_{\rm th}$, the yield of the PQ asymmetry remains a constant given by

$$Y_{\rm PQ} = \epsilon \frac{3T_{\rm th}}{4m_S}. \tag{15}$$

Similarly, with Eq. (2), the angular speed is

$$\dot{\theta}(T) = \epsilon \frac{g_* \pi^2}{30} \frac{T_{\rm th} T^3}{m_S f_{\rm eff}^2(T)}. \tag{16}$$

Based on Eq. (4), we obtain

$$Y_B = \epsilon \frac{3c_B T_{\rm th} T_{\rm ws}^2}{4m_S f_{\rm eff}^2(T_{\rm ws})} \simeq 10^{-10} \epsilon\xi \left(\frac{T_{\rm th}}{100\ {\rm GeV}}\right) \left(\frac{10^9\ {\rm GeV}}{f_{\rm eff}(T_{\rm ws})}\right)^2 \left(\frac{\rm TeV}{m_S}\right). \quad (17)$$

This expression is valid whether thermalization or the electroweak phase transition occurs first and is also general for any type of the potential. While contours of $\epsilon\xi T_{\rm th}$ can be easily included in Fig. 1, a concrete model is necessary to realize the required values of $T_{\rm th}$. In the Supplemental Material [38], we thoroughly demonstrate a consistent thermalization history required by Eq. (17) for the observed baryon asymmetry in the framework of supersymmetry. Supersymmetry is again motivated by the flatness of the potential, or equivalently a light saxion, to obtain a large saxion initial field value. A large viable parameter space is similarly obtained in the supersymmetric version of axiogenesis. In summary, Fig. 1 shows that a wide range of the saxion mass m_S is viable, while a low f_a is favored in the minimal realization of axiogenesis.

Discussion.—We propose a mechanism to explain the baryon asymmetry of the Universe. The two main ingredients are a rotation in the axion direction in the early Universe, corresponding to an excess of PQ charges, as well as QCD and electroweak sphaleron processes that convert the PQ asymmetry into those of baryons and leptons. We construct a concrete model where the rotation is a consequence of higher dimensional PQ-breaking operators. This is analogous to how the rotation of the Affleck-Dine field arises. We show that a sufficient baryon asymmetry is generated from the PQ charge by the QCD and electroweak sphaleron transitions.

Intrinsic to the axiogenesis framework, the angular speed of the rotation needed for the observed baryon asymmetry leads to axion dark matter. In fact, axion dark matter is overproduced in the minimal scenario where the weak anomaly coefficient of the PQ symmetry is as large as the QCD anomaly coefficient and the PQ charge is conserved even after the electroweak sphaleron transition becomes ineffective, which the Standard Model predicts to be at $T_{\rm ws} = 130$ GeV. Therefore, unless the PQ charge is depleted after the electroweak phase transition, the associated prediction is a value of $T_{\rm ws}$ that is higher than predicted by the Standard Model and/or a large weak anomaly coefficient. We show how new physics at the 1–10 TeV scale can raise $T_{\rm ws}$ so that the axion can constitute a subdominant or correct amount of dark matter. In addition to new heavy states, axiogenesis also favors a small decay constant, which is accessible to many axion haloscope and helioscope experiments [72–84]. The evolution of the PQ breaking field reveals nonstandard cosmological eras, which alone may have profound implications for other aspects of cosmology. These phenomenological prospects render axiogenesis an exciting avenue to pursue theoretically and experimentally.

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