Baryogenesis from Strong CP Violation and the QCD Axion

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We show that strong CP violation from the QCD axion can be responsible for the matter antimatter asymmetry of the Universe in the context of cold electroweak baryogenesis if the electroweak phase transition is delayed below the GeV scale. This can occur naturally if the Higgs couples to a $\mathcal{O}(100)$ GeV dilaton, as expected in some models where the Higgs is a pseudo-Nambu-Goldstone boson of a new strongly interacting sector at the TeV scale. The existence of such a second scalar resonance with a mass and properties similar to the Higgs boson will soon be tested at the LHC. In this context, the QCD axion would not only solve the strong CP problem, but also the matter antimatter asymmetry and dark matter.

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Introduction.—Understanding the generation of the matter-antimatter asymmetry of the Universe $\eta = \eta_{10} \times 10^{-10}$, where $\eta_{10} = 6.047 \pm 0.074$ [1], is one of the key motivations for physics beyond the standard model (SM). The SM fails to satisfy two of the three Sakharov conditions needed for baryogenesis. First, as the electroweak (EW) phase transition is not first order [2], there is no sufficient departure from equilibrium in the standard cosmological evolution. Second, CP violation in the SM appears to be too small to explain the value of η. While a large number of possibilities for new sources of CP violation arise in minimal TeV scale extensions of the SM and have been considered for baryogenesis, it is natural to wonder whether the CP nonconserving term in the SM QCD Lagrangian

$$\mathcal{L} = \bar{\Theta}\frac{\alpha_s}{8\pi} G_{\mu\nu a}\tilde{G}_a^{\mu\nu}, \qquad \bar{\Theta} = \Theta + \arg\det \mathbf{M}_q, \tag{1}$$

could have played a role for baryogenesis. The CP-violating $\bar{\Theta}$ term is constrained today to be smaller than 10^{-11} from the absence of a measurable electric dipole moment for the neutron [3]. The Θ parameter characterizes the nontrivial nature of the QCD vacuum which solves the $U(1)_A$ problem [4]. Because chiral transformations change the Θ vacuum once we include weak interactions and the quark mass matrix $\mathbf{M}_q$, the only physical observable angle is $\bar{\Theta} = \Theta + \arg\det \mathbf{M}_q$. Understanding why Θ and arg $\det \mathbf{M}_q$ should be tuned such that $|\bar{\Theta}| < 10^{-11}$ is the so-called strong CP problem [5]. The puzzle is solved if $\bar{\Theta}$ is promoted to a dynamical field which relaxes naturally to zero, as advocated by Peccei-Quinn (PQ) [5]. This solution postulates a new global axial symmetry $U(1)_{\rm PQ}$ spontaneously broken by a scalar field $\Phi = (f_a + \rho(x))e^{ia(x)/f_a}/\sqrt{2}$, where the Goldstone boson $a(x)$ is the axion. New heavy colored quarks with coupling to Φ generate a $G\tilde{G}$ term,

$$\frac{\alpha_s}{8\pi}\frac{a(x)}{f_a} G_{\mu\nu a}\tilde{G}_a^{\mu\nu}. \tag{2}$$

Axion couplings are all suppressed by the factor $1/f_a$ while its mass today satisfies $m_a f_a \approx m_\pi f_\pi$. The axion $a(x)$ relaxes towards the minimum of its potential, at $\langle a \rangle = 0$, thus explaining why $\bar{\Theta}$ is very small today. However, in the early Universe, just after $U(1)_{\rm PQ}$ breaking, $\bar{\Theta} = a(x)/f_a$ is large and frozen to a value of order 1 as long as the axion is massless. The axion acquires a mass at the QCD phase transition and classical oscillations of the axion background field around the minimum of the potential start at T_i when its mass is of the order of the Hubble scale, $m_a(T_i) \sim 3H(T_i) \sim \Lambda_{\rm QCD}^2/M_{\rm Planck}$. The energy stored in these axion oscillations behaves exactly as cold dark matter and is bounded by $\rho_{\rm CDM} \sim 10^{-47}$ GeV$^4 \gtrsim \bar{\Theta}^2 m_a^2 f_a^2/2 \sim \bar{\Theta}^2 m_\pi^2 f_\pi^2$; therefore, today $\bar{\Theta} \lesssim 10^{-21}$. This cosmological constraint leads to an upper bound on f_a whereas an astrophysical lower bound from excess cooling of supernovae, narrows the remaining allowed window for the QCD axion to [6] 10^9 GeV $\lesssim f_a \lesssim 10^{11}$ GeV. The question whether $\bar{\Theta}$ could have played any role at the time of the EW phase transition (EWPT) was investigated only once in the literature in Ref. [7] which concluded that strong CP violation could not be responsible for baryogenesis. We revisit this question here and show that this almost-SM source of CP violation can explain baryogenesis under rather minimal assumptions.

A baryogenesis theory requires a stage of nonequilibrium dynamics in addition to CP violation and baryon number (B) violation. An interesting route is to consider the case where EW symmetry breaking is triggered through a fast tachyonic instability [8]. In this case, the Higgs mass squared is turning negative not as a consequence of the standard cooling of the Universe but because of its coupling to another scalar field which acquires a vacuum expectation value (VEV) and forces the Higgs mass to change rapidly. This "Higgs quenching" leads to the production of unstable EW field configurations [$SU(2)$ textures] which when decaying lead to Chern-Simons number transitions. This dynamics can lead to very efficient production of B at zero temperature and is at the origin of the *cold* baryogenesis scenario [9–11]. Cold baryogenesis requires three conditions: large Higgs quenching to produce Higgs winding number in the first place, unsuppressed CP violation at the time of quenching to bias a net baryon number, and a reheat temperature below the sphaleron freeze-out temperature $T \sim 130$ GeV [12] to avoid washout of B by sphalerons. Sphalerons play no role in this mechanism. Baryogenesis takes place at low temperature, due to the out-of-equilibrium production of SM EW large field configurations. Cold baryogenesis can successfully account for the value of η. It has been simulated on the lattice for different Higgs quenching parameters [11] and using as CP-violating source the dimension-six operator $(\Phi^\dagger\Phi/M^2)\mathrm{Tr}F\tilde{F}$ where Φ is the Higgs doublet, F is the EW field strength and M is constrained to be $M \gtrsim 65$ TeV by the electron electric dipole moment bound. In this Letter, we show that strong CP violation from the operator [Eq. (2)] could source cold baryogenesis instead. Besides, we show that a Higgs-dilaton coupling not only naturally delays the EWPT to temperatures $T \lesssim \Lambda_{QCD}$ but also induces sufficient Higgs quenching during the EWPT, while keeping the reheat temperature below the sphaleron freeze-out temperature, therefore naturally enabling cold baryogenesis. A first step in this direction was provided in Ref. [13].

Baryogenesis from SM baryon number violation.—Because of the EW anomaly, baryogenesis from SM baryon number violation can result from the effective CP-violating coupling

$$\frac{\alpha_W}{8\pi}\zeta(\varphi)\mathrm{Tr}F\tilde{F}, \tag{3}$$

where $\zeta(\varphi)$ is a time-varying function of fields. This generates a chemical potential $\mu \equiv \partial_t\zeta$ for B violation. The resulting B asymmetry is given by

$$n_B = N_F\int dt\frac{\Gamma\mu}{T} \sim N_F\frac{\Gamma(T_{\mathrm{eff}})}{T_{\mathrm{eff}}}\Delta\zeta, \tag{4}$$

where Γ is the rate of Chern-Simons transitions, N_F is the number of families and T_{eff} characterizes the temperature at which Chern-Simons transitions are operative at the same time as efficient CP-violation effects. Using the sphaleron rate in the EW symmetric phase $N_F\Gamma = 30\alpha_w^5T^4 \sim \alpha_w^4T^4$, this leads to

$$\frac{n_B}{s} = \frac{N_F\alpha_w^4 45}{2\pi^2 g_*(T_{\mathrm{reh}})}\left(\frac{T_{\mathrm{eff}}}{T_{\mathrm{reh}}}\right)^3\Delta\zeta \sim 10^{-7}\left(\frac{T_{\mathrm{eff}}}{T_{\mathrm{reh}}}\right)^3\Delta\zeta, \tag{5}$$

where T_{reh} is the reheat temperature after the EWPT and is of the order of the Higgs mass. It may be significantly higher than the temperature of the EWPT, T_{EWPT}, if the EWPT was delayed and completed after a supercooling stage [8]. For standard EW baryogenesis taking place during a first-order EWPT, $T_{\mathrm{eff}} = T_{\mathrm{EWPT}} = T_{\mathrm{reh}}$. In contrast, the key point for cold baryogenesis is that $T_{\mathrm{eff}} \neq T_{\mathrm{EWPT}}$ [8]. T_{eff} should be viewed as an effective temperature associated with the production of low-momentum Higgs modes during quenching. It is significantly higher than T_{EWPT}. It is a way to express the very efficient rate of B violation in terms of the equilibrium expression $\Gamma \sim \alpha_w^4T^4$ although the system is very much out of equilibrium.

Axion-induced CP violation and baryogenesis.—Our goal is to investigate whether the large values of the effective vacuum angle in Eq. (2) at early times can have any implications for EW baryogenesis. We have $\bar{\Theta} = a/f_a \sim \mathcal{O}(1)$ for $T \gtrsim 1$ GeV, and then $\bar{\Theta}$ quickly drops as the axion gets a mass and starts oscillating around the minimum of its potential. The axion Lagrangian reads $\mathcal{L}_a = \mathcal{L}(\partial_\mu a) - \frac{1}{2}\partial^\mu a\partial_\mu a + (a/f_a)(\alpha_s/8\pi)G\tilde{G}$ so that $(\partial V_{\mathrm{eff}}/\partial a) = -(1/f_a)(\alpha_s/8\pi)G\tilde{G}$. Gluon condensation from $SU(3)$ instantons leads to a VEV for $G\tilde{G}$ and a potential for the axion that can be written as $V = f_\pi^2m_\pi^2(1-\cos(a/f_a)) \approx f_a^2m_a^2(1-\cos(a/f_a))$. As a result, $(\alpha_s/8\pi)\langle G\tilde{G}\rangle = f_a^2m_a^2\sin\bar{\Theta}$. To make a connection between the axion and EW baryogenesis, we have to construct an effective operator gathering gluons and EW gauge bosons. An operator of the type [Eq. (3)] can arise, where ζ is controlled by the axion mass squared. In particular, the η' meson, which is a singlet under the approximate $SU(3)$ flavor symmetry of strong interactions, can couple to both $G\tilde{G}$ and $F\tilde{F}$. At temperatures below the η' mass, $m_{\eta'} \approx 958$ MeV, we can use the effective operator

$$\mathcal{L}_{\mathrm{eff}} = \frac{1}{M^4}\frac{\alpha_s}{8\pi}G\tilde{G}\frac{\alpha_w}{8\pi}F\tilde{F}, \tag{6}$$

where $1/M^4 = 10/(F_\pi^2m_{\eta'}^2)$ [7]. We end up with $\mathcal{L}_{\mathrm{eff}} = (1/M^4)\sin\bar{\Theta}m_a^2(T)f_a^2(\alpha_w/8\pi)F\tilde{F}$, hence $\zeta(T) \equiv (1/M^4)\sin\bar{\Theta}m_a^2(T)f_a^2$. The time variation of the axion field and/or mass is thus a source for baryogenesis,

$$n_B \propto \int dt\frac{\Gamma(T)}{T}\frac{d}{dt}[\sin\bar{\Theta}m_a^2(T)]. \tag{7}$$

To estimate the resulting B asymmetry, we will use as the temperature-dependent axion mass with $T_t = 102.892$ MeV [14]: $m_a^2(T) = m_a^2(T=0)$ for $T \leq T_t$ and $m_a^2(T) \approx m_a^2(T=0)\times(T_t/T)^{6.68}$ for $T > T_t$. The axion mass is very suppressed at temperatures above the QCD scale. A large B asymmetry is therefore produced only if the EWPT occurs not much earlier than the QCD phase transition. We have $\Delta\zeta \sim (10 f_a^2/f_\pi^2 m_{\eta'}^2)\delta[\sin\bar{\Theta} m_a^2]|_{\rm EWPT} \sim 0.044 \sin\bar{\Theta}(T_{\rm EWPT}) \times (T_t/T_{\rm EWPT})^{6.68}$ for $T_{\rm EWPT} > T_t$. The final B asymmetry is then, using Eq (5),

$$\frac{n_B}{s} \sim \begin{cases} 4\times 10^{-9}\left(\frac{T_{\rm eff}}{T_{\rm reh}}\right)^3 \sin\bar{\Theta}_{\rm EWPT} \quad \text{for } T_{\rm EWPT} \le T_t, \\ 4\times 10^{-9}\left(\frac{T_{\rm eff}}{T_{\rm reh}}\right)^3 \sin\bar{\Theta}_{\rm EWPT}\left(\frac{T_t}{T_{\rm EWPT}}\right)^{6.68}, \\ \text{for } T_{\rm EWPT} > T_t. \end{cases} \tag{8}$$

In the context of standard EW baryogenesis, we have $T_{\rm eff} = T_{\rm EWPT} = T_{\rm reh}$. Even if we set $\sin\bar{\Theta}(T_{\rm EWPT}) \sim 1$ in Eq. (8), we need $T_{\rm EWPT} \lesssim 0.2$ GeV, to get a large enough B asymmetry today. Such a low EWPT temperature can be achieved by coupling the Higgs to a dilaton field [15,16] whose scalar potential energy induces a supercooling stage. However, the reheat temperature $T_{\rm reh} \sim \mathcal{O}(m_H, m_d)$ cannot be kept below a GeV unless the dilaton with mass m_d is very weakly coupled. Besides, in this case, the axion oscillations would be delayed and would overclose the Universe. This led Kuzmin *et al.* to conclude that strong CP violation from the axion cannot play any role during EW baryogenesis [7]. We conjecture, on the contrary, that the axion can well explain baryogenesis. To deduce the B asymmetry, we include in Eq. (8) the temperature dependence of $\bar{\Theta}$ in terms of $\bar{\Theta}_i$, the initial value when oscillations start at T_i defined by $3H = m_a(T_i)$. The energy density stored in axion oscillations redshifts as nonrelativistic matter, leading to $\bar{\Theta}^2(T) = \bar{\Theta}_i^2(m_a(T_i)/m_a(T))(T/T_i)^3$ for $T \le T_i$. For temperatures above the QCD phase transition, the B asymmetry is suppressed by the axion mass, while at low temperatures, it is suppressed by the smallness of $\bar{\Theta}$ if the axion started to oscillate in the supercooling stage, i.e., $m_a \gtrsim 3H_{\rm EW} \sim \sqrt{3\rho_{\rm vac}}/m_{\rm Pl}$. In that case, it is then typically maximized for EWPT temperatures in the 10 MeV–1 GeV range, as shown in the left Fig. 1. The gray curve, which corresponds to the standard EW baryogenesis assumption, $T_{\rm eff} = T_{\rm reh}$, reproduces the negative conclusion of Ref. [7]. At the time, the mechanism of cold baryogenesis was not known. On the other hand, cold baryogenesis cures the problem. The key point is that even if $T_{\rm EWPT} \lesssim \Lambda_{\rm QCD}$, we can have $T_{\rm eff} \gtrsim T_{\rm reh} \sim m_H$. From lattice simulations of cold baryogenesis [11], a quenched EWPT typically has $(T_{\rm eff}/T_{\rm reh}) \sim 20$–$30$. Considering that there is no washout factor, we easily get a large B asymmetry since $n_B/s \propto (T_{\rm eff}/T_{\rm reh})^3$. For $m_a \lesssim 3H_{EW}$ (right plot), there is no low-T suppression as the axion field value does not start decreasing until after the EWPT so that a large B asymmetry can be produced even for a very delayed EWPT. We indicated the dependence on the initial angle $\bar{\Theta}_i$ by a band corresponding to the range $10^{-2} \le \bar{\Theta}_i \le \pi/2$.

Axion dark matter.—If $f_a \lesssim 7\times 10^{10}$ GeV, axion oscillations start during supercooling and the corresponding energy density gets diluted compared to the usual prediction. For larger f_a, axion oscillations do not start until after reheating thus the contribution to DM from the axion misalignment mechanism is unchanged compared to the standard cosmological evolution. On the other hand, the contribution to DM from the decay of strings is reduced by at least $\sim \log(\Lambda_{\rm EW}/\Lambda_{\rm QCD}) \sim 7$ e folds of inflation associated with the supercooling stage. Together with the B asymmetry, this is the only other relevant observable consequence of our supercooling stage. Our peculiar cosmological scenario has no impact on big bang nucleosynthesis since the Universe is reheated well above the MeV scale. Besides, there is no constraint from electric dipole moment bounds since CP violation from $\bar{\Theta}$ is large only in the early Universe while it is very suppressed today

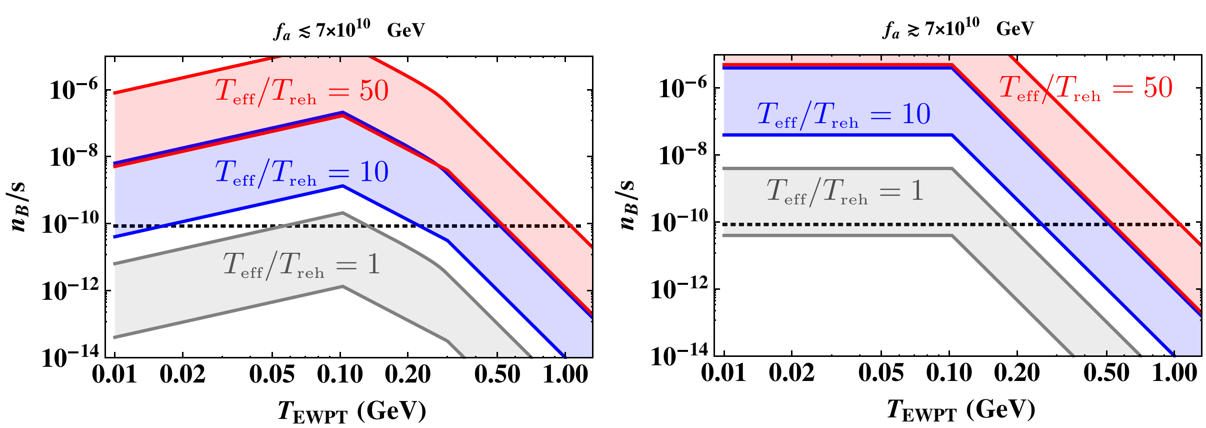

FIG. 1 (color online). Today's baryon asymmetry as a function of $T_{\rm EWPT}$ compared with measured value (dotted line). The case $T_{\rm eff}/T_{\rm reh} = 1$ and $T_{\rm eff} \sim T_{\rm EWPT}$ that would characterize standard EW baryogenesis is unfeasible as $T_{\rm reh} \sim \mathcal{O}(m_H) \gg \Lambda_{\rm QCD}$. The cases with $T_{\rm eff}/T_{\rm reh} \gtrsim 10$ can easily account for a large B asymmetry and correspond to a quenched EWPT, as in cold EW baryogenesis. Each band corresponds to varying the initial angle $\bar{\Theta}_i$ in the range $[10^{-2}, \pi/2]$. Left: $m_a \gtrsim 3H_{\rm EW} \sim 3\times 10^{-14}$ GeV, oscillations start at $T = 0.3$ GeV in the supercooling era before the EWPT. Right: $m_a \lesssim 3H_{\rm EW}$, the axion is frozen to its initial value until after reheating.

after the axion has relaxed to its minimum. We conclude that the standard QCD axion can be responsible for both dark matter and the B asymmetry of the Universe in the context of cold EW baryogenesis. We now review the conditions for successful cold EW baryogenesis.

Higgs quench from a Higgs-scalar coupling.—The key point in this Letter is to exploit the fact that efficient B violation can take place at temperatures below the sphaleron freeze-out temperature, under strong out-of-equilibrium conditions as provided by a quenched EWPT. In the standard picture of cold baryogenesis [8,11], the tachyonic transition develops when the Higgs mass squared m_{eff}^2 changes sign rapidly due to a coupling of the Higgs to an additional scalar field. Just before the EWPT, the Universe is relatively cold. The dynamics of spinodal decomposition has been investigated both analytically and numerically [10,17–21], typically using infinitely fast quench. The Fourier modes of the Higgs field with low momentum are unstable and grow exponentially. These extended field configurations play a key role in inducing Chern-Simons transitions (see, e.g., [13] for a summarized review and references therein). Although we are far from thermal equilibrium, the rate of Chern-Simons transitions can be approximated by that of a system in thermal equilibrium at a temperature T_{eff}. The production of nonzero Chern-Simons number when the Higgs field experiences a fast quench depends on the Higgs mass and on the speed of the quench. For $m_H \sim 125$ GeV and quenching parameter $|u| \gtrsim 0.1$, where $u \equiv (1/2m_H^3)(dm_{\mathrm{eff}}^2/dt)|_{T=T_q}$, lattice simulations [11] found $T_{\mathrm{eff}}/T_{\mathrm{reh}} \sim$ 20–30. In the SM, the effective Higgs mass varies solely because of the cooling of the Universe so that $u^{\mathrm{SM}} \sim (1/\mu^3)(d/dt)(-\mu^2 + cT^2)|_{T=T_q} \sim (T_{\mathrm{EW}}/M_{\mathrm{Pl}}) \sim 10^{-16}$. This situation changes radically if the Higgs mass is controlled by the time-varying VEV of an additional field σ, e.g., $m_{\mathrm{eff}}^2(t) = \mu^2 - \lambda_{\sigma\phi}\sigma^2(t)$ which leads to $u = -(\lambda_{\sigma\phi}/m_H^3)(\sigma\dot{\sigma})|_{t_q}$. From energy conservation $(\dot{\sigma})^2 \sim \mathcal{O}(\Delta V)$, and we can naturally get order 1 quenching parameter as it is no longer controlled by the Hubble parameter. This additional coupling of the Higgs is what the cold baryogenesis scenario assumes. In this Letter, we provide a natural motivation for such an assumption, building up on Ref. [13].

Naturally delayed EW phase transition and low reheat temperature from the dilaton.—We consider the following scalar potential in which the quadratic term for the Higgs field ϕ is controlled by the VEV of the dilaton σ,

$$V(\sigma, \phi) = V_\sigma(\sigma) + \frac{\lambda}{4}(\phi^2 - \xi\sigma^2)^2, \tag{9}$$

where $V_\sigma(\sigma)$ is a scale invariant function modulated by a slow evolution $V_\sigma(\sigma) = \sigma^4 \times P(\sigma^\epsilon)$, where $|\epsilon| \ll 1$. ϕ is the Higgs field and $\xi = v^2/f^2$ is a constant. Note that this potential is precisely the one of Randall-Sundrum models [22], where $\sigma \equiv ke^{-k\pi r}$ is the radion field. The scale $f \equiv \langle\sigma\rangle$ is generated once the radion is stabilized. The cosmology of the potential [Eq. (9)], $V_\sigma(\sigma)$, was summarized in Ref. [16]. The value of the field at tunneling, σ_r, is $\sigma_r \sim \sqrt{\sigma_+\sigma_-}$, where σ_+ and $\sigma_- = f \sim \mathcal{O}(\mathrm{TeV})$ are the positions of the maximum and minimum of the potential, respectively. The nucleation temperature is $T_n \sim 0.1\sigma_r \sim 0.1f\sqrt{\sigma_+/\sigma_-}$. For a standard polynomial potential, $\sigma_+ \sim \sigma_- \sim \sigma_r \sim T_n$. In contrast, for the very shallow dilatonlike potential, $\sigma_+ \ll \sigma_-$, and the nucleation temperature is parametrically much smaller than the scale associated with the minimum of the potential. We therefore naturally get a stage of supercooling before the phase transition completes. The hierarchy between σ_- and σ_+ can be as large as the Planck scale or weak scale hierarchy, $\sigma_-/\sigma_+ \lesssim \Lambda_{\mathrm{UV}}/f$, and T_n can then be as low as $T_n \sim 0.1f\sqrt{f/\Lambda_{\mathrm{UV}}}$ [16]. We obtain $T_n \sim 35$ MeV if $f = 5$ TeV and $\Lambda_{\mathrm{UV}} = M_{\mathrm{Pl}}$, while $T_n \sim 0.1$ GeV if $f = 1$ TeV and $\Lambda_{\mathrm{UV}} = f_{\mathrm{PQ}} = 10^{10}$ GeV. While a delayed EWPT down to the QCD scale is a general outcome in our framework, the breaking of conformal invariance by QCD will modify the scalar potential V_σ around the QCD scale and affect detailed predictions.

In the cases considered so far for cold baryogenesis, the quenching time is defined when the effective Higgs mass vanishes, which translates as $\sigma_q^2 = \mu^2/\lambda_{\sigma\phi}$. For the dilaton-like potential [Eq. (9)], $\mu = 0$, and quenching should happen between tunneling and the time the field rolls down to the minimum. The condition $u \gtrsim 0.1$ becomes $(\xi\lambda/m_H^3)\sqrt{2(V_\sigma(0) - V_\sigma(\sigma))}\sigma \gtrsim 0.1$. Depending on the value of λ, the quenching condition can be realized relatively well before the field reaches the minimum. When estimated at the minimum, it translates as a bound on the second scalar mass eigenstate, $(\lambda v^2 m_d/m_H^3) \sim (m_d/2m_H) \gtrsim 0.1$, where $m_d = \sqrt{V_\sigma''(\sigma = f)}$. Since we are considering $\xi = v^2/f^2 \ll 1$, the off diagonal terms in the squared mass matrix of the Higgs-dilaton system are small compared to the diagonal entries. Therefore, the two mass eigenvalues are essentially $2\lambda v^2$ and $V_\sigma''(f)$.

Having checked that the quenching criterion is readily fulfilled, the last condition to be satisfied for successful cold baryogenesis is that the reheat temperature does not exceed the sphaleron freeze-out temperature to prevent washout of the asymmetry. After the EWPT, the vacuum energy stored in the Higgs and dilaton fields reheats the plasma, $(8\pi g_* T_{\mathrm{reh}}^4/30) = \Delta V$ where $\Delta V \sim m_d^2 f^2$. Imposing $T_{\mathrm{reh}} < 130$ GeV leads to a constraint on the dilaton mass. For $f \sim \mathcal{O}(\mathrm{TeV})$, this means that m_d should be $\mathcal{O}(100)$ GeV. Constructions that lead naturally to such a light dilaton have been recently discussed in Ref. [23–27]. LHC constraints on an EW scale dilaton were presented before the Higgs discovery in [28–31]. Interpretation of the Higgs discovery in terms of a Higgs-like dilaton [32] has then been considered in [33,34]. We are instead interested in a scenario where in addition to the 125 GeV Higgs, there is a light dilaton, which is a less constrained option, see, e.g., [35–37], and a careful analysis of CMS and ATLAS data is generally definitely worthwhile and will be a key test for our scenario in particular.

Conclusion.—The QCD axion could play a key role in providing the new source of *CP* violation in baryogenesis, therefore linking the origin of dark matter to that of the matter antimatter asymmetry of the Universe. Because of a Higgs-dilaton coupling, the EW phase transition is naturally delayed to sub-GeV temperatures. The subsequent reheating stage can dilute the axion particles produced by string decays but does not modify the usual prediction from the axion oscillations which start only well after reheating. It would be interesting to refine these dark matter predictions which depend on the details of the reheating process. The experimental tests of this scenario are of three very different types: the usual QCD axion searches, LHC searches for an additional dilatonlike field coupled to the Higgs, and a stochastic millihertz gravity wave background [16] detectable by eLISA.

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