QCD-Electroweak First-Order Phase Transition in a Supercooled Universe

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If the electroweak sector of the standard model is described by classically conformal dynamics, the early Universe evolution can be substantially altered. It is already known that—contrarily to the standard model case—a first-order electroweak phase transition may occur. Here we show that, depending on the model parameters, a dramatically different scenario may happen: A first-order, six massless quark QCD phase transition occurs first, which then triggers the electroweak symmetry breaking. We derive the necessary conditions for this dynamics to occur, using the specific example of the classically conformal $\mathcal{B}$-$\mathcal{L}$ model. In particular, relatively light weakly coupled particles are predicted, with implications for collider searches. This scenario is also potentially rich in cosmological consequences, such as renewed possibilities for electroweak baryogenesis, altered dark matter production, and gravitational wave production, as we briefly comment upon.

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Introduction.—Despite the recent discovery of the Higgs boson h, we still have little clue on the physics beyond the standard model (SM) at or above the electroweak (EW) scale. In the past decades, model building has been mostly focusing on supersymmetric or Higgs compositeness scenarios at the (sub-)TeV scale, motivated by the naturalness of the Higgs mass value. These approaches are, however, under strain, due to tighter and tighter experimental bounds on the masses of new particles, notably of colored ones, predicted in such models. Hence, there is renewed motivation to explore alternatives, notably theories including very weakly coupled particles, possibly lighter than SM ones.

An old theoretically appealing idea is that EW symmetry breaking (EWSB) is induced by radiative corrections to the Higgs potential, manifesting conformal symmetry at tree level if its mass term vanishes [1]. This possibility is nowadays excluded in the SM due to the measured values of its parameters, but it might be viable in classically conformal (CC) extensions of the SM where at least an additional scalar field ϕ is introduced, a requirement anyway needed to account for neutrino masses [2] at least if originating through a seesaw mechanism. In a frequently considered implementation of this scenario [3–5,7], it has been noted that the phase transition (PT) breaking the EW symmetry tends to be strongly first order. Then a significant supercooling below the critical temperature and a relatively long time scale for bubble percolation are implied, and thus a sizable gravitational wave (GW) production and possibly electroweak baryogenesis (EWBG) [8–14] are expected. See also [15–17] for other realizations with hidden strong dynamics.

Despite their conceptual simplicity, CC models may lead to an even more fascinating possibility: If the supercooling is maintained down to temperatures lower than the QCD critical temperature $T_c^{\rm QCD}$, chiral symmetry breaking (χSB) occurs spontaneously via quark condensation, $\langle\bar{q}q\rangle \neq 0$. Contrarily to the current phase of the Universe, all the quarks were then massless, as initially the scalar fields have no vacuum expectation value (VEV). The chiral symmetry is thus broken from $SU(6)_L \times SU(6)_R$ to $SU(6)$, and the associated PT is then first order [18]. At the same time, $\langle\bar{q}q\rangle \neq 0$ also breaks the EW symmetry, since $\bar{q}q$ is an $SU(2)_L$ doublet with a nonvanishing $U(1)_Y$ charge, a situation that has been recently considered as relevant only in "gedanken worlds" [19]. Furthermore, when χSB occurs, the Yukawa couplings y_i with the SM Higgs h generate a linear term $y_i h\langle\bar{q}_i q_i\rangle/\sqrt{2}$, tilting the scalar potential along the h direction. This tilt destabilizes the false vacuum at the origin, and the Higgs acquires a VEV *at the QCD scale*. A similar possibility within the SM had been already entertained by Witten [20] but is long since excluded. A couple of decades ago, it was occasionally reconsidered in SM extensions with a dilaton field [21] or in applications to EWBG [22]. The goal of this Letter is to show that it currently remains a concrete possibility in CC extensions of the SM, implying a qualitatively different history of the early Universe. In the following, we focus on characterizing the conditions for a QCD-induced EWSB, commenting upon some particle physics and cosmological consequences of such a scenario.

The model.—For definiteness, let us consider the CC $\mathcal{B}$-$\mathcal{L}$ extension of the SM [4], where the $\mathcal{B}$-$\mathcal{L}$ symmetry is gauged, with gauge coupling g. Besides the SM particles, the model contains a gauge boson Z', a scalar Φ with $U(1)_{\mathcal{B}\text{-}\mathcal{L}}$ charge 2, and three right-handed neutrinos (RHν) canceling the $[U(1)_{\mathcal{B}\text{-}\mathcal{L}}]^3$ and gravitational anomalies. The

CC assumption requires that the scalar potential V, within renormalizable field theories, involving Φ and the Higgs doublet H has no quadratic terms and is given, up to a constant term, by $V(H,\Phi) = \lambda_h|H|^4 + \lambda_{\rm mix}|H|^2|\Phi|^2 + \lambda_\phi|\Phi|^4$. It is then assumed that the $\mathcal{B}$-$\mathcal{L}$ symmetry is radiatively broken by the Coleman-Weinberg (CW) mechanism [1], which triggers the EWSB. For this, the scalar mixing $\lambda_{\rm mix}$ is required to take a small negative value. Here we summarize approximate key formulas [5]. At one loop, the CW potential along the potential valley is approximated by $V_{\rm CW}^v(\phi) = V_0 + B\phi^4[\ln(|\phi|/M) - 1/4]/4$, where $\phi \approx \sqrt{2}|\Phi|$. If the condition $B > 0$ is satisfied, the effective potential has a global minimum at $\phi = M$, the scale generated radiatively via the RG evolution of the quartic coupling (see [4]). Note that the zero temperature one-loop CW potential $V_{\rm CW}^v$ has no barrier between $\phi = 0$ and M. At the global minimum, various particles acquire masses: $m_{Z'} = 2gM$, $m_\phi = \sqrt{B}M$, and $m_{N_i} = Y_i M/\sqrt{2}$ for the RHν's whose Yukawa couplings are Y_i. The constant term $V_0 = BM^4/16$ is chosen so that $V_{\rm CW}^v(M) = 0$. The coefficient B is approximately given by $B \approx c_0^4[3(2g)^4 - D\lambda_{\rm mix}^2 - 2{\rm Tr}(Y/\sqrt{2})^4]/8\pi^2$, where Tr is the trace over three RHν flavors, $c_0^4 \approx (1 + \lambda_{\rm mix}/\lambda_h)$, and $D \simeq 41$ (see Supplemental Material [5]). The coefficient c_0^4 and $D \neq -1$ represents the admixture of h and ϕ along the valley. Because of $\lambda_{\rm mix} < 0$, the Higgs field $h = \sqrt{2}|H|$ has a minimum at $v = (|\lambda_{\rm mix}|/2\lambda_H)^{1/2}M$, identified with the Higgs VEV 246 GeV. If $M \gg v$, the Higgs mass is given by $m_h \approx \sqrt{|\lambda_{\rm mix}|}M$. In spite of the large hierarchy $M \gg v$, the scalar ϕ is generically light, $m_\phi \lesssim 10 m_h$ [23]. In the following, for simplicity we shall require $g \gtrsim 10^{-4}(m_{Z'}/{\rm TeV})^{1/4}$ so that ϕ and Z' are thermalized before the supercooling stage, although some of our conclusions may be true in a broader parameter space.

Hypercooling in the EW sector.—A CC system has peculiar thermodynamic properties. To see this, let us focus on a model accounting only for the fields ϕ and Z' and at the leading order in the high-temperature expansion (see Supplemental Material [5] for some considerations on the quality of these approximations). The effective potential is thus approximated as (see, e.g., [24]) $V(\phi) = (c_2/2)T^2\phi^2 - (c_3/3)T\phi^3 + (B_{Z'}/4)\phi^4\ln(T/\hat{\mu})$, where ϕ-independent terms are dropped. The coefficients are given by $c_2 = g^2$, $c_3 = 6g^3/\pi$, $B_{Z'} = 6g^4/\pi^2$, and $\hat{\mu} = m_{Z'}e^{\gamma_E - 1/2}/4\pi$. At sufficiently high T, the quadratic term dominates and the only minimum of the potential is at $\phi = 0$: $\mathcal{B}$-$\mathcal{L}$ symmetry is restored. We study the cosmological evolution of the Universe with an initial condition $\phi = 0$, which is naturally realized after the large field inflation as discussed, e.g., in Refs. [25–27]. When T drops below the critical temperature $T_c \sim m_{Z'}$, defined by the condition $\mathcal{C}(T_c) \equiv 9c_2B\ln(T_c/\hat{\mu})/(2c_3^2) = 1$, the nontrivial minimum of the potential at $\phi_c = 3Tc_2/c_3 \lesssim M$ has a lower energy compared to the false vacuum $\phi = 0$. However, due to the CC assumption, the coefficient of the quadratic term c_2 is always positive, and the false vacuum remains the local minimum even at $T \ll T_c$: The thermal potential barrier never disappears. Hence, the Universe with the Hubble expansion rate $\mathcal{H} = \sqrt{V_0/3m_{\rm pl}^2}$ is supercooled down to a very low temperature where $m_{\rm pl} \approx 2.4 \times 10^{18} \times$ GeV is the reduced Planck mass.

The Universe may eventually percolate into the true vacuum via bubbles nucleated by quantum tunneling. The percolation temperature T_p can be estimated by using the tunneling rate $\Gamma \approx T^4 e^{-S_3/T}$. In the present model, the critical bubble's action is given by $S_3/T \approx A[1 - 2\pi\mathcal{C}(T)/9]^{-1}$, where $A = 43.7c_2^{3/2}/c_3^2 \propto g^{-3}$ and $\mathcal{C}(T) = (3/4)\ln(T/\hat{\mu})$ for $T < \hat{\mu}$, consistently with results in Ref. [28] for a non-negative quartic coupling. Thus, for $g \ll 1$ the tunneling rate becomes very small. The fraction of space remaining in the false vacuum at a given temperature $T < T_c$ is given by $p(T) = e^{-I(T)}$, where $I(T)$ is defined by the probability that a single bubble of true vacuum is nucleated in the past (see [29]). The percolation temperature T_p is then defined by the condition $I(T_p) = 1$. In Fig. 1, we plot T_p as a function of g and $m_{Z'}$. Because of the (weakly T-dependent) behavior $S_3/T \propto g^{-3}$, percolation does not occur for $g \lesssim 0.2$. Eventually, the transition to the true vacuum would occur when the de Sitter fluctuation $\mathcal{O}(1) \times T_{\rm GH}$ becomes comparable to the width of the barrier, $\Delta\phi \sim T/g$, where $T_{\rm GH} \equiv \mathcal{H}/2\pi$ is the Gibbons-Hawking temperature. This does not happen until the temperature becomes very low when de Sitter fluctuation destabilizes the false vacuum, a condition that we dub *hypercooling*. Note that the width of the barrier is evaluated as $\Delta\phi \sim T/g$ and implies that $m_{Z',{\rm eff}} < g\Delta\phi \sim T$, where the high-$T$ expansion is barely justified. Our conclusion remains qualitatively correct in a more realistic treatment, e.g., using the full thermal potential without high-T expansion [30].

QCD-induced EWSB.—If the percolation temperature of the $\mathcal{B}$-$\mathcal{L}$ sector is lower than the QCD critical temperature $T_c^{\rm QCD}$ and if the de Sitter fluctuation $\sim T_{\rm GH}$ is negligible

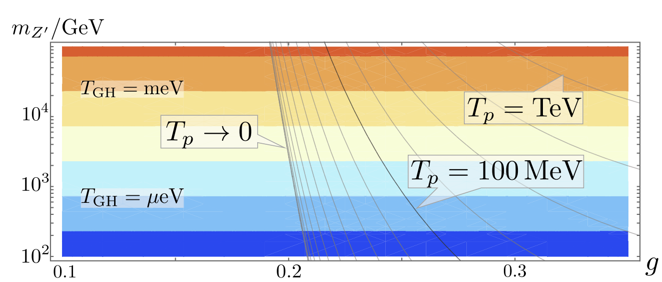

FIG. 1. Contour plot of the percolation temperature T_p (black lines) as a function of g and $m_{Z'}$. The horizontal color bands show the temperature $T_{\rm GH} \equiv \mathcal{H}/2\pi$.

compared to the QCD scale, the previous model cannot be trusted anymore to describe the dynamics, since the CC condition is actually broken by QCD via dimensional transmutation, i.e., confinement and χSB. At the false vacuum, all the quarks are massless. QCD with $N_f = 6$ massless quarks (or five massless and one massive near the false vacuum) has a first-order PT [18], with $T_c^{\rm QCD}$ somewhat lower than that in the SM, e.g., 85 MeV in Ref. [33]. Contrarily to the previously discussed case, the QCD PT is expected to occur at $T_n^{\rm QCD}$ only mildly below $T_c^{\rm QCD}$, because QCD has a dynamical scale $\Lambda_{\rm QCD}$. We can check that hypercooling does not take place, e.g., by using the Polyakov-quark-meson model [34].

When the QCD PT occurs, namely, when the chiral condensates form, a linear term $\sum_i y_i \langle \bar{q}_i q_i \rangle h/\sqrt{2}$ is generated in the Higgs potential, and a new local minimum $h = v_{\rm QCD} \sim \mathcal{O}(100)$ MeV emerges. At this minimum, quarks (even the top quark) acquire very light masses $m_{q_i} = y_i v_{\rm QCD}/\sqrt{2} \lesssim \Lambda_{\rm QCD}$. Thus, all the $N_f = 6$ quarks are expected to form a chiral condensate $\langle \bar{q}_i q_i \rangle$. The top Yukawa coupling y_t sets the size of the linear term in the Higgs potential; i.e., the local minimum of the Higgs potential is estimated as $v_{\rm QCD} = (y_t \langle \bar{t}t \rangle / \sqrt{2}\lambda_h)^{1/3}$. Note that the top behaves similarly to the strange quark in the present Universe, which has a mass $m_s \sim 100$ MeV comparable with the QCD scale, but whose condensate is of the same order as the up (or down) quark one.

Also, the $SU(2)_L \times U(1)_Y$ gauge symmetries are spontaneously broken, and linear combinations of the pions and the ordinary Nambu-Goldstone components of the Higgs field are eaten by the massive gauge bosons. Thus, EWSB is triggered by the first-order QCD PT.

Histories of the early Universe.—Different histories of the early Universe, i.e., different trajectories of the scalar fields, are possible as in Fig. 2, depending on different values of the parameters $(g, m_{Z'})$ as in Fig. 3. If the percolation temperature T_p is higher than the QCD scale $\Lambda_{\rm QCD} \sim 100$ MeV, ϕ field tunnels into the true vacuum before the QCD PT (green line in Fig. 2). A strong first-order PT takes place, and a sizable production of gravitational waves is expected [10]. From Fig. 1, such a possibility is realized for sufficiently strong gauge coupling $g \gtrsim 0.2$ as in the green region in Fig. 3.

If $g \lesssim 0.2$, the QCD-induced EWSB occurs after de Sitter expansion with an e-folding $\sim \ln(T_i/T_n^{\rm QCD})$. $T_i \equiv (30V_0/\xi\pi^2)^{1/4}$ is the temperature when the expansion starts, and $\xi \gtrsim 110$ is number of degrees of freedom in the extended SM. After the QCD-induced EWSB, if the quadratic term $(c_2 T^2 + \lambda_{\rm mix} h^2/2)\phi^2/2$ at $T = T_n^{\rm QCD}$ and $(\phi, h) = (0, v_{\rm QCD})$ is positive, the field ϕ is trapped [dubbed scenario (II)]. If it is negative, ϕ rolls down to the true minimum $(\phi, h) = (M, v)$, which we name scenario (I). The trajectories are drawn in Fig. 2. The trapping condition translates into $6m_{Z'}^2 + {\rm Tr}(m_N^2) > 12 m_h^2 (v_{\rm QCD}/T_n^{\rm QCD})^2$ and is reported in pink in Fig. 3. Then thermal inflation occurs: As T drops, the field tunnels and starts rolling around $T = \sqrt{|\lambda_{\rm mix}|/(2c_2)}\, v_{\rm QCD}$ when the coefficient of the quadratic term vanishes. If, instead, the trapping condition is violated, ϕ freely rolls down [37] [scenario (I), yellow region in Fig. 3]. On top of it, the fate of the Universe is also controlled by the slow roll condition $|\eta| = m_{\rm pl}^2 |V''|/V_0 < 1$ at $(0, v_{\rm QCD})$. Namely, if $g \lesssim 10^{-2}(m_{Z'}/{\rm PeV})^3$ is satisfied, an inflationary expansion takes place after the phase transition.

Since the CW mechanism requires $B > 0$, i.e., $3m_{Z'}^4 > 2{\rm Tr}(m_N^4) + Dm_h^4$, the necessary condition for scenario (I) reads $v_{\rm QCD}/T_n^{\rm QCD} > [D/12]^{1/4} \simeq 1.36$, i.e., $\langle \bar{t}t \rangle^{1/3} \gtrsim 0.77 T_n^{\rm QCD}$. In the standard two-flavor QCD case, $\langle \bar{q}q \rangle^{1/3}/T_n^{\rm QCD} \simeq 0.5$, and one falls a little short of this condition. However, pending dedicated lattice studies, in our framework we cannot exclude that this inequality is actually satisfied. Anyway, as long as one is near the condition $3m_{Z'}^4 \gtrsim Dm_h^4$, either scenario (I) is realized or scenario (II) with very shallow trapping, i.e., very short inflation and a fast transition to the true vacuum. This parameter space provides ideal conditions to observe, e.g., relics from the QCD-induced EWSB, limiting dilution. It is

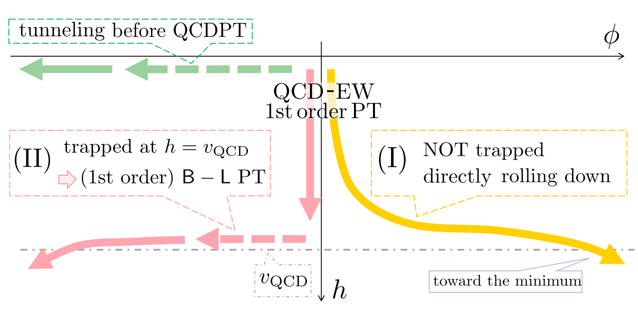

FIG. 2. Possible trajectories of the scalar fields (ϕ, h) in the early Universe. All start from the origin (0,0).

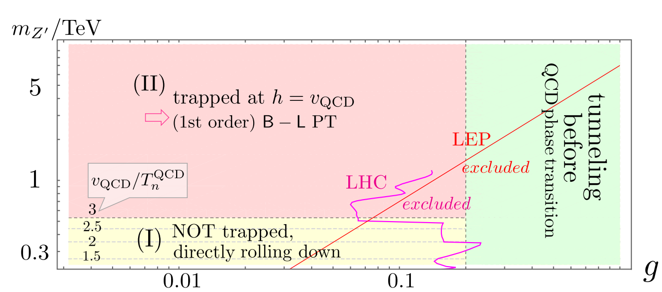

FIG. 3. Schematic cosmological histories for the parameter region $m_{Z'} \sim$ TeV, assuming $m_N = 0$ and $v_{\rm QCD}/T_n^{\rm QCD} = 3$. The horizontal thin dashed lines are the boundaries between (I) and (II) for $v_{\rm QCD}/T_n^{\rm QCD} = 1.5$, 2, and 2.5. Below the lower edge, $B > 0$ is violated. For reference, we plot the large electron positron bound [35] (red line) and some LHC bounds from Fig. 6 in Ref. [36] (magenta line).

also the regime where the $\mathcal{B}$-$\mathcal{L}$ gauge boson is predicted at the EW scale and RHν's and the $\mathcal{B}$-$\mathcal{L}$ scalar below it, which makes the model amenable to direct collider probes like the LHC and, up to ~4 GeV mass for RHν, also SHiP: See, for instance, [39] for a forecast study, reporting a sensitivity down to $g \sim \mathcal{O}(10^{-3})$.

Cosmological consequences.—We list a few cosmological consequences of the above scenario.

(i) The temperature after the PT is limited to $T_i \lesssim$ 20 GeV in scenario (I). Hence, particles with mass $m \gtrsim \mathcal{O}(10) \times T_i$ (such as many dark matter candidates) cannot be thermally produced. The viability of different types of dark matter candidates obtained via alternative production mechanisms should be thus revisited (see, e.g., [17]).

(ii) *Cold* EWBG might take place, which has been argued to be a generic opportunity offered by the supercooling stage ending with the first-order PT [40]. An interesting possibility is a QCD axion extension [41]. As discussed in the *standard* EWBG context [22], the EWPT triggered by the χSB "optimizes" the efficiency of the strong *CP* violation to that purpose. Of course, our scenario is very specific, and a modification of the EWSB dynamics has profound implications on several ingredients of the EWBG scenario, like the sphaleron energy and the necessary *CP* violation.

(iii) Note that the *e*-folding $\ln(T_i/T_n^{\rm QCD})$ gained during the late inflationary period is small and unrelated to the one probed via cosmic microwave background (CMB) fluctuations. This is a welcome consequence of our model, since small-field inflations with simple symmetry-breaking potential (including the CW one) are otherwise inconsistent with observations. Models like Ref. [42], where the CW inflation with a Higgs linear term comes from the quark condensate, should be reanalyzed within the present framework.

(iv) Another consequence of the first-order QCD PT is that the formation of primordial black holes (PBHs; see, e.g., [43]) as well as of primordial magnetic fields [44] is eased. If PBHs form, due to the horizon size at the QCD scale their mass is predicted in the (tenths of) solar mass range: They might contribute to the dark matter of the Universe, can be searched for via lensing, and, being massive enough, through accretion disks may alter the heating and ionization history between CMB recombination and first star formation, with consequences for CMB observables as well as for future 21 cm probes [45].

(v) The most direct cosmological probe would consist in the detection of the GW background produced via bubble collisions. Following the standard formulas [46,47], the GW power spectrum is determined by $\beta/\mathcal{H}$, where $\mathcal{H}$ is the Hubble parameter at the production of GWs and $1/\beta$ corresponds to the duration of the PT and the typical size of the bubbles at the collision. The parameter $\beta/\mathcal{H}$ is hard to compute reliably, although it is expected to be larger than ~100 under reasonable assumptions [48]. An additional parameter is $\zeta \equiv T_{\rm rh}/T_i$, with $T_{\rm rh}$ the reheating temperature, which quantifies the duration of the reheating period where the scalar oscillates around the true minimum behaving like pressureless matter. In Fig. 4, we illustrate the approximate GW signal expected under the assumption $T_i = 10$ GeV with varying ζ and $\beta/\mathcal{H}$. It is worth stressing that, in scenario (I), $\beta/\mathcal{H}$ is essentially independent from g, in contrast to $\beta/\mathcal{H} \propto g^{-2}$ in scenario (II) [49]. For comparison, the sensitivities of three configurations foreseen for the space mission LISA [50] are also reported.

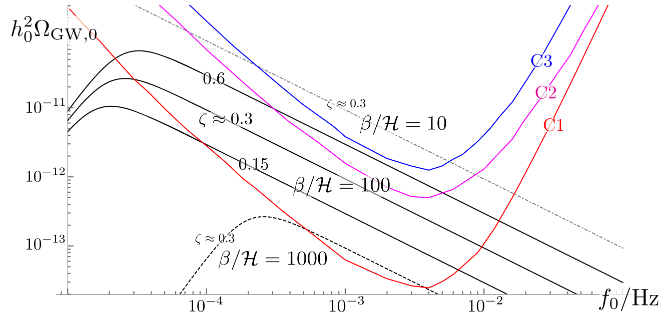

FIG. 4. The GW power spectrum for $\beta/\mathcal{H} = 10, 100, 1000$. We chose $T_{\rm i} = 10$ GeV and $\zeta = 0.6, 0.3, 0.15$. The sensitivity of three configurations foreseen for the space mission LISA [50] are also shown.

Conclusions.—Phenomenologically, we know only that the thermal history of the Universe is conventional below temperatures of a few MeV, sufficient to set up the initial conditions (e.g., populating active neutrino species) for primordial nucleosynthesis. It is usually assumed that the knowledge of the SM allows one to backtrack the evolution of the Universe up to temperatures of few hundreds of GeV and that the EWPT is a crossover, as predicted by the SM, although theories with an extended EW sector where a first-order EWPT occurs are not rare. It is, however, almost universally accepted that the QCD PT is not first order, *even in models of physics beyond SM*, hence with very limited implications for the later Universe. Here we offer a counterexample, where an extension of the SM motivated only by the EW physics sector changes both the QCD and EW PT dynamics, with the possibility of a very peculiar history of the Universe: A first-order QCD PT (with six massless quarks) triggers a first-order EWPT, eventually followed by a low-scale reheating of the Universe where hadrons (likely) deconfine again, before a final, "conventional" crossover QCD transition to the current vacuum. To the best of our knowledge, this is the only viable scenario known where a first-order QCD PT can be obtained without large lepton [51] or baryon asymmetry [52]. We have only sketched some important particle physics and cosmological consequences of this scenario. The actual reach of forthcoming collider searches, the extension to more general models than the $\mathcal{B}$-$\mathcal{L}$ here used for illustration, as well as quantitative consequences for cosmological crucial problems such as dark matter or baryon asymmetry are all interesting aspects which we plan to return to in the near future.

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Note added.—Recently, we became aware of an ongoing, related study [53] motivated in the context of Randall-Sundrum models.

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