Quantum chromodynamic preheating: new frontier of baryogenesis

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The dynamic time-dependent vacuum of quantum chromodynamics (QCD) provides a nonperturbative production mechanism of cosmological matter abundance beyond the thermodynamic framework. The mechanism is similar to the so-called preheating induced by the nonadiabatic-varying vacuum. In the proposed scenario, the dynamic light quark condensate acts as if it is an inflaton field of the conventional preheating. This is a new type of particle production scenario via the nonperturbative QCD preheating, which is fulfilled before the quark condensate reaches the true vacuum and the static hadron phase is formed in the evolving Universe. The baryon and antibaryon pairs created by the QCD preheating are overproduced by two orders of magnitude more than those estimated from the thermal abundance. Baryogenesis is possible once the asymmetry \( \epsilon = (n_{N} - \bar{n}_{N})/(n_{N} + \bar{n}_{N}) \) of \( O(10^{-9}) \) is provided from a new sector external to QCD, where \( n_{N} \) and \( \bar{n}_{N} \) denote the baryon and antibaryon number densities, respectively. Thus the dynamic aspect of the QCD vacuum opens a new frontier to explore low-scale baryogenesis around the subatomic scale.

PACS numbers:

Introduction – The nucleon mass is supplied almost fully by the nonperturbative nature of QCD, which amounts to about ninety-nine percent of the full mass. The dynamical mass generation is thought to be responsible for the spontaneous breaking of the approximate chiral symmetry at the nonperturbative vacuum. The rest, about a few percent, comes from a couple of the chiral breaking sources arising externally to QCD: small masses of up and down quarks (called the current quark masses) originated from the Higgs in the electroweak theory, and the electromagnetic interaction for the charged up and down quarks. This is a well-known feature associated with the static aspect of the QCD vacuum in the chiral broken and confining phase. Then, what could the dynamic aspect of the QCD vacuum provide for the nucleon?

The dynamic aspect of vacuum in quantum field theory has extensively been explored so far in light of the reheating process after the inflation in the early Universe, that is called the preheating [1-5] (for reviews, see e.g., [5][8]). The oscillating background field of the inflaton induces the nonadiabatic state for other coupled species. It drives the nonperturbative particle production, and the amount is exponentially amplified by the parametric resonance similar to a swing with a pumping oscillator. The mechanism of preheating is applied to interesting variant scenarios: the production of massive particles heavier than the inflaton [6][12]; the preheating due to the alternative field instead of the inflaton [13][16], etc.

The “ballpark” of preheating has also been extended to the baryogenesis via the nonperturbative production at high scales [17][28]. However, preheating at the subatomic scale has never been discussed in the context of baryogenesis. A scalar condensate is present even in QCD, that is the light quark condensate \( \langle \bar{q}q \rangle \), and it can couple to the nucleon state as well as meson states in a systematic way respecting the chiral symmetry and its breaking by the current quark masses. Therefore, the dynamic motion of \( \langle \bar{q}q \rangle \) should have the potential to explosively produce the number densities for nucleon and anti-nucleon by nonadiabatic processes, similarly to the preheating induced by the nonadiabatic-varying vacuum.

In this Letter, we prove that the observation above is indeed correct: the QCD preheating takes place by the dynamic \( \langle \bar{q}q \rangle \), and the dynamic aspect of the QCD vacuum provides the production mechanism of the baryon and antibaryon pairs. Thus the main production mechanism of the cosmological matter abundance is present in QCD, and is provided by the same source as the generation of the main part of the nucleon mass. See Fig. 1.

We monitor the dynamic QCD vacuum by working on a low-energy chiral effective theory where the role of the QCD “inflaton” field \( \langle \bar{q}q \rangle \) is replaced by the interpolating mesonic field \( \sim \sigma \). The couplings of \( \sigma \) to the low-lying hadron spectra including the nucleon are fixed by the chiral symmetric and breaking structures reflecting what the underlying QCD possesses.

Since the CP and baryon number violations in the standard model are too small to explain the observed amount, we introduce an additional dark sector communicating with a part of the QCD sector. This setup is similar to the aforementioned introduction of the external-chiral...
Another QCD baryogenesis has been addressed based on higher
break a source arising from the electroweak theory, so as to gain the realistic nucleon mass.

We find that the QCD preheating predicts overproduction of the baryon and antibaryon number densities, which are by two orders of magnitude greater than those estimated from the thermal equilibrium distribution at around the QCD phase transition, as described in the text, where the QCD preheating via the dynamic \( \langle \sigma \rangle \sim \langle \bar{q}q \rangle \) happens at \( T \sim 0.7T_{pc} \sim 109 \text{ MeV} \) by starting the roll-down in the potential, and reheats the Universe up to \( T = T_{reh} \simeq 117 \text{ MeV} \) (see the text). The bottom panel traces the time evolution of \( T \) in terms of the potential shapes of the dynamic \( \langle \bar{q}q \rangle \).

The thermal corrections during the QCD preheating are negligible because the time scale of the preheating is much shorter than the relaxation time scale, as clarified in Eq. (10) in the later section. Passing the relaxation, the system goes back to the normal thermal equilibrium, which makes the potential of \( \langle \sigma \rangle \) lifted up in the standard manner of thermal QCD.

**Overview of the QCD preheating scenario** – To begin with, we outline what we know about the static aspect of the QCD/chiral phase transition, and promote what we can suspect for the dynamic feature of the QCD vacuum.

The lattice QCD has confirmed that thermal QCD with \( 2 + 1 \) flavors at the physical point undergoes the crossover for the chiral phase transition, which is called the chiral crossover, at the pseudocritical temperature \( T_{pc} \sim 155 \text{ MeV} \) [25–29]. At almost the same temperature, the deconfinement-confinement transition (crossover) is expected to happen as well [30, 31]. Above \( T_{pc} \), the light quark condensate \( \langle \bar{q}q \rangle \) takes nearly vanishing values when properly renormalized to be divergent free. Cooling down to \( T_{pc} \), \( \langle \bar{q}q \rangle \) starts to get sizable and finally reaches the value measured at the vacuum when \( T \) goes well below \( T_{pc} \), \( T/T_{pc} \lesssim 0.7 \) [25].

It is in a single and static volume that the lattice calculation has observed this crossover. In contrast, the Universe was expanding and evolving even during the QCD phase transition, which should be described by the formation of “bubbles” and their evolution. The barrier-less boundary of a single “bubble” between the chiral symmetric/quark-gluon and broken/hadron phases is thought to be given as crossover which starts to happen at \( T = T_{pc} \), that is consistent with what the lattice simulation has observed. The hadron-phase bubbles expand until all bubbles occupy the Universe, which time corresponds to the epoch where \( \langle \bar{q}q \rangle \) reaches the true vacuum (at \( T \lesssim 0.7T_{pc} \)). This implies that the evolution of the chiral broken/hadron phase acts as if the Universe undergoes a supercooling in the regime \( 0.7T_{pc} \lesssim T \lesssim T_{pc} \) [2]. In the present Letter, we apply this supercooling-like picture to the dynamic \( \langle \bar{q}q \rangle \) during the chiral crossover epoch (from \( T_{pc} \) down to \( 0.7T_{pc} \)) [2], in such a way that \( \langle \bar{q}q \rangle \) keeps staying at \( \sim 0 \), i.e., \( \langle \bar{q}q \rangle \sim 0 \), with the vanishing velocity, until the completion of the chiral broken/hadron phase at \( T \lesssim 0.7T_{pc} \).

The dynamic motion of \( \langle \bar{q}q \rangle \) will globally be processed instantaneously at \( T \sim 0.7T_{pc} \). Since the scalar field \( \langle \bar{q}q \rangle \) couples to the meson and baryon states, the nonadiabatic particle production will then take place and reheat the Universe from 0.7T_{pc} up to \( T_{reh} \lesssim T_{pc} \). The typical time scale of the dynamic \( \langle \bar{q}q \rangle \) motion is much faster than the background Hubble evolution, as will be clarified later.

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\#1 Another QCD baryogenesis has been addressed based on higher scale features of QCD, not at the subatomic scale [34], where the production mechanism of the number densities is neither due to the nonperturbative time-dependent varying vacuum.

\#2 In the literature [33, 34], a similar observation has been made and the barrier-less “bubble” formation reflecting the chiral crossover was discussed based on the Gibbs energy conservation between the quark-gluon and hadron phases, which is applied to a context different from the present QCD preheating.

\#3 The supercooling could create spatial inhomogeneity of the Universe. The presently discussed supercooling epoch for \( 0.7T_{pc} < T < T_{pc} \) is compatible with the discussion in [35] on the allowed amount of inhomogeneity in the QCD-phase transition epoch.
Hence the $T$ dependence is safely negligible until the $\langle \bar{q}q \rangle$ motion ends to be trapped at the true vacuum. This is the characteristic feature of the QCD preheating, in contrast to a class of the preheating scenarios at higher scales (see, e.g., a review [3–8]), where the Hubble friction becomes crucial during the inflaton oscillation. This is how the QCD preheating works and can successfully create the pions and nucleons through the nonperturbative and nonadiabatic processes, as will be explicitly seen below.

A schematic view of the QCD preheating is illustrated in Fig. 1.

**Dynamic QCD vacuum via a chiral effective theory description** – The dynamic $\langle \bar{q}q \rangle$ across the chiral crossover at around $T = T_{pc}$ communicates both the chiral symmetric and broken phases of QCD. In accordance with this picture, in the present work we employ a chiral effective theory described by hadronic-interpolating fields for the color-singlet composite operators $\bar{q}q$ and $\bar{q}q_q$, that is conventionally called the linear sigma model.

The model is based on the chiral $SU(2)_L \times SU(2)_R$ symmetry with the current quark mass for the up and down quarks included. The building blocks are: i) the complex two by two scalar matrix: $M = M_{\pi \pi} + M_{\pi N}$, which transforms under the chiral symmetry as $M \rightarrow U^\dagger_{L,R} \cdot M \cdot U^\dagger_{L,R}$ with $U_{L,R} \in SU(2)_{L,R}$, where $M$ is parametrized by the isosinglet sigma ($\sigma$) mode and isrotplet pion ($\pi$) mode as $M = \sigma \cdot 2_{\times 2} / 2 + i \pi^a \tau^a / 2$ with the Pauli matrices $\tau^a$ ($a = 1, 2, 3$); ii) the nucleon (proton, neutron)-doublet field $N_{L,R} = (p_n)^{\dagger}_{L,R}$, which belong to the fundamental representation of $SU(2)_{L,R}$ groups. The linear-sigma model Lagrangian is thus given as

$$
\mathcal{L} = \text{tr} \left[ \partial_{\mu} M^\dagger \partial^\mu M \right] - V + \bar{N}i\partial T N - 2m_N f_{\pi} \left( \bar{N}_L M N_R + \bar{N}_R M^\dagger N_L \right),
$$

$$V = m^2_{\sigma} f_{\pi} \text{tr} \left[ \text{Re}(M) \right] + m^2 \text{tr} \left[ M^\dagger M \right] + \lambda \left( \text{tr} \left[ M^\dagger M \right] \right)^2.
$$

$f_{\pi}$ is the pion decay constant $\simeq 92.4$ MeV, and $m_N$ the nucleon mass $\simeq 940$ MeV. All the terms are invariant under the chiral symmetry, C and P symmetries, except for the first term coupled to the pion mass $m_\pi \simeq 140$ MeV, which breaks the chiral symmetry in a way reflecting the explicit quark mass term in the underlying QCD. The explicit-chiral or -isospin breaking for the nucleon sector has been neglected for simplicity. The light quark condensate generated in QCD is aligned to be two-flavor diagonal (isospin symmetric) and P invariant, due to the Vafa-Witten’s theorem [30], hence the dynamics of the vacuum expectation value of $\sigma$, $\langle \sigma \rangle$, monitors the dynamic $\langle \bar{q}q \rangle$.

We fix the potential parameters $m^2$ and $\lambda$ to the values determined at the true vacuum $\langle \sigma \rangle = f_{\pi}$ satisfying the stationary condition. Then we have $\lambda f_{\pi}^2 = (M_{\sigma}^2 - m_\sigma^2) / 2$ and $m^2 = 3m_\pi^2 / 2 - M_{\sigma}^2 / 2$, with the $\sigma$ mass squared defined as $M_{\sigma}^2 = \partial^2 V / \partial \sigma^2 |_{\sigma = f_{\pi}}$, which we take

$\simeq (500 \text{ MeV})^2$.

The equation of motion for $\langle \sigma \rangle$ with the space homogeneity leads

$$0 = \langle \dot{\sigma} \rangle + \gamma \langle \dot{\sigma} \rangle - m_\pi^2 f_{\pi} + m^2 (\langle \sigma \rangle + \lambda \langle \sigma \rangle^3 + \cdots). \quad (2)
$$

The ellipses denote the negligible terms including the Hubble friction term $\langle 3H \langle \dot{\sigma} \rangle \rangle$ and the backreactions from the pion and nucleon fields.

We have added another friction term with the coefficient $\gamma$, which plays the role of the full width of the $\sigma$ meson when it is identified as $f_{0}(500)$ in the Particle Data Group. As a phenomenological input, we take $\gamma$ to be the central value of the current measurement, $\gamma \equiv |2 \text{Im} \sqrt{\langle \sigma \rangle_{\text{pol}}}| = 550$ MeV [27]. This phenomenological $\gamma$ fully includes nonperturbative scattering contributions, so we can discard the interaction terms with the pions because those should fully be included in $\gamma$. Note also that compared to this $\gamma$ term, the Hubble friction term is now safely negligible because $H \sim T^2 / M_{pl} \ll \gamma$ with $M_{pl}$ the Planck mass $M_{pl} \sim 10^{19}$ GeV.

The backreaction terms correspond to the effective mass of $\sigma$ arising as the plasma effect through the nonadiabatic production of the pion and the nucleons due to the dynamic $\sigma$. We have checked that their effect on the motion of $\langle \sigma \rangle$ is sufficiently small. The nonadiabatic particle production occurs around $\langle \sigma \rangle \sim 0$ at once, but the parametric resonance does not since the later motion does not come back there due to the friction $\gamma$. Thus the dynamic $\langle \sigma \rangle$ is well approximately described only by Eq. (2) without interaction terms with pions and nucleons. The time evolution of $\langle \sigma \rangle$ is plotted in Fig. 2 as a function of $(M_{pl} t)$.
Since the dynamic $\langle \sigma \rangle$ rolls down from $\langle \sigma \rangle \sim 0$ to $f_\pi$ with the strong instantaneous, the potential energy $V(\sigma = 0) - V(\sigma = f_\pi) = \frac{1}{2}(\mathcal{M}_\pi^2 + 3m_N^2)f_\pi^2 \approx (135 \text{ MeV})^4$ is converted into the radiation energy. The reheating temperature can thus be estimated as

$$T_{\text{reh}} = \left[4\langle T_{\text{pc}} \rangle^4 + \frac{30}{\pi^2}g_*(135 \text{ MeV})^4\right]^{1/4} \approx 117 \text{ MeV},$$

where we have assumed the relativistic degrees of freedom $g_* \approx 20$ at $T(\langle \sigma \rangle = 0) = 0.7T_{\text{pc}} \approx 109 \text{ MeV}$, in which the dynamic $\langle \sigma \rangle$ starts to roll. The produced entropy density can also be estimated as

$$s = \frac{2\pi^2}{45}g_* \cdot (117 \text{ MeV})^3 \approx 1.41 \times 10^7 \text{ MeV}^3.$$

### Nonadiabatic production of nucleons due to the dynamic $\langle \sigma \rangle$ – The dynamic $\langle \sigma \rangle$ controls the mass of the nucleons through the Yukawa interaction in Eq. (1):

$$\frac{2m_N}{f_\pi} \vec{N}_L M N_R = \frac{m_N}{f_\pi} \langle \sigma \rangle \cdot \vec{N}_L N_R + \cdots .$$

Therefore, the nucleon mass $\tilde{m}_N(t) = m_N \frac{\langle \sigma(t) \rangle}{f_\pi}$ varies in time, following the time evolution of $\langle \sigma \rangle$ depicted in Fig. 2. Such a time-varying mass causes the nonperturbative nucleon production during the violation condition of the adiabaticity, $|\tilde{m}_N/\tilde{m}_N^2| \gtrsim 1$. This inequality leads to the production area as

$$|\sigma| \lesssim \sqrt{\frac{f_\pi}{m_N} \langle \sigma \rangle} \approx 42 \text{ MeV},$$

where we estimated $\langle \sigma \rangle \approx \sqrt{V(\sigma = 0) - V(\sigma = f_\pi)} \sim (135 \text{ MeV})^2$. Hence, the nonperturbative nucleon production would be completed within the region in Eq. (6) that in terms of Fig. 2 corresponds to

$$M_\rho t \lesssim 5 .$$

The actual production time can be earlier than the number in Eq. (7) because the estimated velocity $\langle \dot{\sigma} \rangle$ would be smaller due to the friction $\gamma$.

The total nucleon number density can be evaluated as

$$n_N(t) + \bar{n}_N(t) = \sum_k \frac{\langle \dot{\rho}_N(\vec{k}, t) \rangle}{\omega(\vec{k}, t)} ,$$

for $N = p, n$. Here $\dot{\rho}_N(\vec{k}, t)$ is the kinetic energy density of the nucleon in momentum space, which is derived from the Hamiltonian as

$$\dot{\rho}_N(\vec{k}, t) = \frac{1}{V} \left[ \vec{N}_L(\vec{k}, L) \left( \vec{\gamma} \cdot \vec{k} + \tilde{m}_N(t) \right) N_L(\vec{k}, t) \right. + \left. \vec{N}_R(\vec{k}, t) \left( \vec{\gamma} \cdot \vec{k} + \tilde{m}_N(t) \right) N_R(\vec{k}, t) \right] + 2\omega(\vec{k}, t) ,$$

with $\omega(\vec{k}, t) = \sqrt{\vec{k}^2 + \tilde{\tilde{m}}^2(t)}$ being the one-particle energy of the nucleon and $V = \int d^3x$ the space volume of the system and $N_L(\vec{k}, t)$ and $N_R(\vec{k}, t)$ the Fourier transformed Dirac fields. The last term of (9) corresponds to the subtraction of the negative vacuum energy ($4 \times \frac{1}{2} \omega(\vec{k}, t)$). The time evolution of the total baryon number density in Eq. (8) is thus determined by solving coupled equations of motion for $N_L(\vec{k}, t)$, $N_R(\vec{k}, t)$, and $\langle \sigma(t) \rangle$, which is plotted in Fig. 2.

We find the total number is explosively generated by the nonadiabatic-time varying $\langle \sigma(t) \rangle$, in the time range $M_\rho t \lesssim 3$, and gets asymptotically saturated to be $(n_N + \bar{n}_N) \approx 10^3 \text{ MeV}^3$.

Note that this number density is much larger than the thermal equilibrium density at the reheating temperature $T_{\text{reh}} \approx 117 \text{ MeV}$, $[(n_N + \bar{n}_N)]_{\text{EQ}} \approx 3 \times 10^3 \text{ MeV}^3$, hence becomes overabundant. Actually, the overproduced nucleons can survive long enough during the whole reheating process: the relaxation time scale, at which the overproduced nucleons pair-annihilate, can be estimated as $M_\rho \Delta t_{\text{relax}} = M_\rho (n_N(\sigma v))^{-1} \sim 700$, or equivalently

$$\Delta t_{\text{relax}} \sim 5 \times 10^{-7} \text{ fs} ,$$

where the static nucleon-pair annihilation cross section has been evaluated as a classical disc $\langle \sigma v \rangle \sim 4\pi/m_N^2$ with the impact parameter $(1/m_N)$. Thus the QCD preheating is operative to stock a large number of nucleon and antinucleon pairs in out-of-equilibrium until the relaxation time. This fact tempts us to consider the application to the baryogenesis, in which the Sakharov criteria are required: the baryon number violating interaction, C- and CP-violating interactions, and out-of-equilibrium condition [35].

Once the baryonic asymmetry $\epsilon \equiv (n_N - \bar{n}_N)/(n_N + \bar{n}_N)$ is provided by the other mechanism until the relaxation time, the QCD preheating can successfully generate the desired amount of the net baryon number of the Universe,

$$Y_B \equiv \frac{(n_N + \bar{n}_N)}{s} \cdot \epsilon \sim 10^{-10} \times \left( \frac{\epsilon}{10^{-9}} \right) ,$$

where the entropy density in Eq. (4) is applied. Here a salient feature is seen: the observed baryon number can be realized by a relatively smaller asymmetry than that accumulated in the thermal equilibrium due to the overproduced nucleons, which will not be washed out as long as the asymmetry generation completes before the system goes back to the thermal equilibrium.

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#4 $\vec{A}$ denotes the spatial component of the vector variable $A$. 
It is true that QCD in the standard model cannot solely generate the size of the asymmetry $\epsilon$ in Eq. (11) because of the absence of sufficient CP- and baryon number-violating interactions, but Eq. (11) opens a new roadmap as new baryogenesis involving a number of the new sector candidates transferred to the standard model: once an external sector transferring the CP and baryon number violations to QCD is hypothesized, the QCD preheating triggered by the dynamic $\langle \bar{q}q \rangle$ can create the net baryon asymmetry through Eq. (11), if and only if it ends until the relaxation time $M_\sigma \Delta t_{\text{relax}} \sim 700$. This is our main result.

A source of asymmetry – Successful realization of the baryogenesis scenario calls for an additional sector that includes the CP- and baryon number-violating sources and supplies $\epsilon$ in Eq. (11). There may be various candidate models for a scenario of this sort. In the present work, we introduce a class of models in which the baryogenesis makes the most of the QCD preheating.

Consider a dark sector with dark-Majorana neutron fields $n_{D_L,R}$ being allowed to couple to the QCD neutron fields $n_{L,R}$ in a minimal way:

$$
\mathcal{L}_{n-n_D} = -m_D \bar{n}_D n_D - \frac{1}{2} M_L \bar{n}_L n_L - \frac{1}{2} M_R \bar{n}_R n_R - g_L \bar{n}_L n_D - g_R \bar{n}_R n_L + \text{h.c.},
$$

where the superscript $c$ stands for the charge conjugation. In general, only two CP-violating phases can be physical to be introduced among the mass couplings above. For simplicity, we turn off one of two phases and embed it in $M_L$. With this phase, the nonzero Majorana mass couplings thus transfer the CP violation and the baryon number violation into the QCD sector.

The dark baryons need to be as heavy as the neutron, in such a way to make the dark-sector communication with the neutron operative in the particle production process.

In the dark sector scenario described in Eq. (12), the dynamic $\langle \sigma(t) \rangle$ also plays an important role in generating the CP-violating source. Indeed, even the physical CP phases on the mass parameters can be erased by the appropriate diagonalization of the mass matrix since the Lagrangian consists of only the two-point interactions. However, the neutron's mass currently depends on time via $\langle \sigma(t) \rangle$ as in Eq. (5), and thus the CP phases can reappear from the kinetic terms as the so-called Berry connection. A similar discussion can be seen in [10]. The combined effect of the CP phase and the baryon number violating couplings thus enables the simultaneous asymmetric production of the neutron and antineutron by the dynamic $\langle \sigma \rangle$.

Production of baryon asymmetry – We shall evaluate the net baryon number instead of the asymmetry $\epsilon$ #6, which can be defined by the (approximately conserved) $U(1)$ Noether charge as

$$
n_B(t) = \frac{1}{V} \int d^3 x \left( \langle n_L^\dagger n_L \rangle + \langle n_R^\dagger n_R \rangle - \sum_k 2 \right),
$$

where $V = \int d^3 x$ is the space volume, and the last term corresponds to the subtraction of the divergent part induced by the zero-point energy $(4 \times \frac{1}{2})$. The time evolution of the vacuum expectation values in Eq. (13) is evaluated by solving chained equations of motion for left- and right-handed neutron fields coupled to the dark left- and right-handed neutrons through Eq. (12), together with the dynamic $\langle \sigma(t) \rangle$ obeying Eq. (2).

The definition of the net baryon number in Eq. (13) should potentially include the proton contribution as well, which, however, we can safely ignore in our analysis. This is because the proton does not have the CP- and baryon number-violating interactions, and thus the net baryon number by the proton cannot be generated.

The sigma motion makes the net baryon number produced through the baryon-number violating couplings $g_{L,R}$. We have observed that the net number starts to oscillate even after the baryon number production. This is because of the presence of the $n-n_D$ and $n-\bar{n}$ oscillations that last eternally. To obtain the static net baryon number, the couplings $g_{L,R}$ connecting the neutrons and the dark sector need to somehow get damping in time or vanishing at the later era. The desired function form would be like a damping oscillation, such as $g_{L,R}(t) = \varphi(t) g_{L,R,0}$ with $\varphi(t) \sim e^{-t \Gamma} \cos m_\varphi t$ #7 for $\Gamma_\varphi, m_\varphi \ll M_\sigma$. In that case, since the magnitude of the $g_{L,R}$ couplings asymptotically and promptly drops to zero, any phenomenological and astrophysical bounds can safely be satisfied when the observation is made.

Figure 3 shows the time evolution of the created asymmetry $Y_B = n_B/s$, normalized by the entropy density

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#6 In the present dynamic system, it is actually less costly to work on the evaluation of the dynamical net nucleon number instead of a direct estimate of $\epsilon$.

#7 The damping-oscillation factor $\varphi(t)$ could arise when one considers an underlying picture of the $g_{L,R}$ mass-mixing coupling to be given by a scalar or dilaton, which communicates between the dark sector and the normal QCD sector. In that case, $\varphi(t)$ would be regarded as a background part of the scalar field, and the form of the damping oscillation in time would be provided by the scalar decay width $\Gamma_\varphi$, following the time evolution equation: $\langle \varphi \rangle + \Gamma_\varphi \langle \varphi \rangle + m_\varphi^2 \langle \varphi \rangle = 0$. This ultraviolet completion could be rich in phenomenology and cosmology, which is, however, beyond the current scope, to be pursued elsewhere.
The details of the numerical computations are to be presented in another publication.

The reference values of the time scales are specific mainly to the currently selected scalar mass $m_\varphi = 10$ MeV. However, the general trend of $Y_B$ as viewed in Fig. 3 is qualitatively independent of $m_\varphi$: larger $m_\varphi$ merely leads to the later second and third phases.
