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Baryogenesis from leptomesons

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Various new physics models, e.g., theories of compositeness, can accommodate the color singlet excited leptons that interact with the leptons, quarks, leptoquarks, etc. A particular type of excited lepton, which can be produced through the four-fermion interaction with a lepton and quark-antiquark (or lepton-antilepton) pair, we call a leptomeson. These new particles may contribute to a variety of experimental anomalies such as the discrepancy in the muon $g-2$. We propose that the leptomesons also can generate the baryon asymmetry that explains the imbalance in ordinary matter and antimatter in the observable Universe. We consider the two types of scenarios for this baryogenesis via leptogenesis to occur from either leptomeson oscillations or decays. Neither possibility contradicts the small masses of the observable neutrinos and the proton stability. Moreover, they can be relevant for the near-future collider experiments and do not suffer from the gravitino problem.

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I. INTRODUCTION

The Standard Model (SM) in particle physics is in good agreement with the majority of the experimental data. However, it does not explain some fundamental issues, e.g., the large number of "elementary" fermions and their arbitrary masses and mixings, the fractional electric charge of the quarks, the similarity between the leptons and the quarks (analogous three flavors and similar behavior under the $SU(2)_{\mathrm{L}} \times U(1)_{\mathrm{Y}}$ symmetry with the same weak coupling), etc. The models of compositeness [1–8] try to solve these problems by introducing the substructure of the SM particles. Theories with a colored substructure of the leptons besides frequently discussed $SU(3)_{\mathrm{c}}$ triplet leptoquarks (LQs) and octet leptogluons may include also $SU(3)_{\mathrm{c}}$ singlet excited leptons, which have larger masses, but same lepton numbers as the SM leptons. So the leptons can be "excited" to these new heavy states by the interactions with other SM particles. A particular type of the excited leptons presents a hypothetical fermion that effectively couples to lepton and the pair of the SM fermion and antifermion. This coupling conserves the baryon number (B) and does not spoil the stability of the proton. We refer to the excited lepton of this type as the leptomeson (LM).[1] In particular, LMs may have the same preon content as lepton-meson pairs.

One example of LM generation can be given in the haplon models [7,11], which are based on the symmetry $SU(3)_c \times U(1)_{em} \times SU(N)_h$ and contain the two categories of colored preons (haplons): the fermions $\alpha^{-1/2}$ and $\beta^{+1/2}$ and the scalars $x^{-1/6}, y^{+1/2}, \ldots$. In this framework, the preon pairs can compose the SM particles as $\nu = (\bar{\alpha}\bar{y})_1$, $d = (\bar{\beta}\bar{x})_3$, $W^- = (\bar{\alpha}\beta)_1$, etc., and the new heavy composites, e.g., LQ $(\bar{x}y)_{\bar{3}}$ and leptogluon $(\bar{\beta}\bar{y})_8$, where the subindex indicates $SU(3)_c$ representation. However, there also can exist multipreon LM states such as $\bar{\beta}\bar{x}\bar{y}x$, $\bar{\alpha}\bar{y}\bar{\beta}\bar{x}\beta x$, etc. This possibility gets more points from recent discoveries of the multiquark states [12,13] due to the similarity between QCD and haplon dynamics. Essentially, LMs can be lighter than the LQs and the leptogluons due to the absence of color dressing. Notice that the possible contribution to the muon $g-2$ from a particular type of LM, which can couple to a lepton and a meson, was discussed in Ref. [14].

One of the most important observations, which cannot be explained within the big bang cosmology and the SM, is the baryon asymmetry (η_B) of the Universe that appears to be populated exclusively with baryonic matter rather than antibaryonic matter [15]. Possible scenarios of dynamical generation of η_B during the evolution of the Universe from a hot early matter-antimatter symmetric stage are known as the baryogenesis (BG) mechanisms. Majority of these scenarios discussed in the literature satisfy the three Sakharov conditions [16]:

(i) Violation of B symmetry
(ii) Violation of C and CP symmetries (to produce an excess of baryons over antibaryons)
(iii) A departure from thermal equilibrium (since the average of B is zero in equilibrium).

Some "exotic" mechanisms of BG that do not satisfy at least one of these conditions were discussed in Refs. [17–20].

The SM does not provide a successful BG due to the lack of CP violation, and not strongly first-order electroweak phase transition (PT) [21] to achieve the departure from thermal equilibrium. However, in the economical SM extensions, η_B can be generated through the thermal

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[1]Notice that in Refs. [9,10] the same term "leptomeson" was used for the bound states of colored excitations of e^+ and e^-.

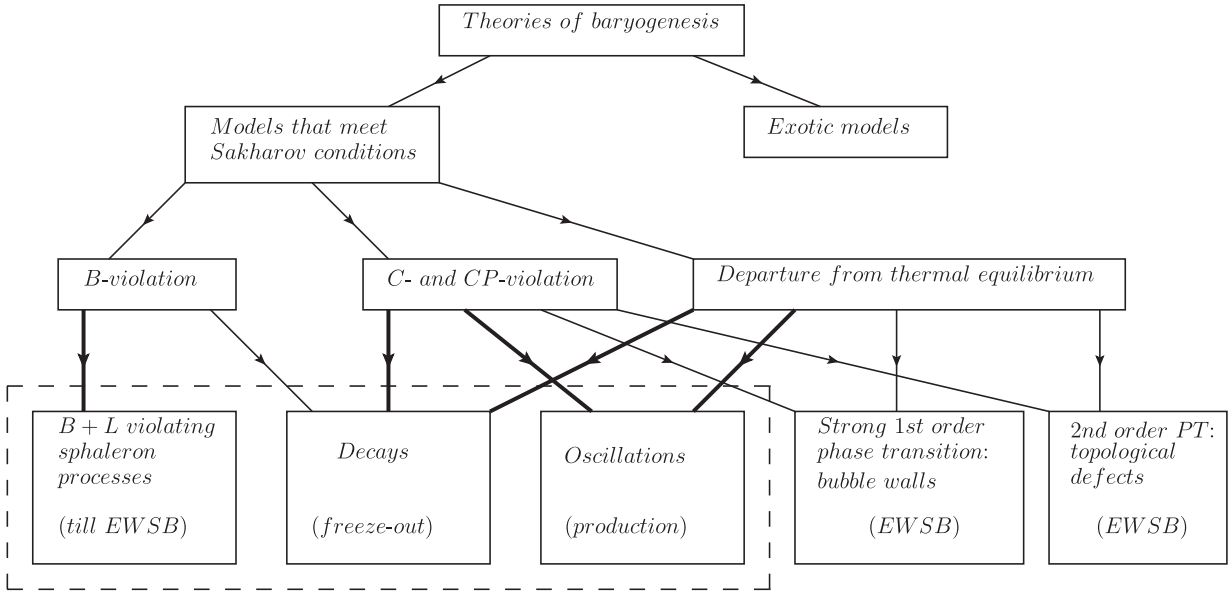

FIG. 1. Structural scheme for various types of baryogenesis and ways to meet Sakharov conditions.

leptogenesis (LG) mechanism [22,23] where the L asymmetry is produced in the out-of-equilibrium decays of heavy Majorana particles and, further, the SM sphaleron processes [24,25] convert this lepton asymmetry into the baryon one. These sphaleron transitions are effective until the electroweak symmetry breaking (EWSB).

However, LG in the supersymmetric generalizations of the SM suffers from the gravitino problem [26–28], which comes from the too high reheating temperature related to the strong lower bound on the right-handed neutrino mass (Davidson-Ibarra bound) [29–32]. To avoid this problem, the resonant mechanisms of LG were introduced [33–39].

In this paper, we investigate how LMs may provide successful BG. The deviation from thermal equilibrium can occur during production (so-called BG from oscillations) [40,41] as well as during freeze-out and decay [22].[2] Depending on the properties of LMs, either one of these two scenarios can be realized in nature. The former case can work for both Dirac and Majorana LM masses and is of particular interest since it can be successful with the LM masses of order of the EWSB scale that can be tested nowadays. The later case requires Majorana masses of LMs similar to the standard LG [22,42–45] from the $SU(2)_{\mathrm{L}}$ singlet neutrino decays. However, the important difference is that the Davidson-Ibarra bound on the heavy neutrino masses, which comes from their see-saw connection to the light neutrino masses through the Yukawa couplings, is not applicable to the considered LM masses. As a result, the LM masses are allowed to be much smaller than the permitted heavy neutrino mass scale of $M_N \gtrsim 10^9$ GeV in the standard LG.

In the flowchart for the BG models shown in Fig. 1 the relevant to present consideration ways to satisfy the Sakharov conditions are emphasized by the bold arrows and the related blocks are encircled by the dashed line. Notice that the models that satisfy these conditions in a non-typical way such as in Ref. [46] are not specified in this flowchart.

In the sections II and III we present possible BG mechanisms from LM oscillations and decays, respectively. We discuss the issue of neutrino masses and conclude in the section IV.

II. BARYOGENESIS FROM LEPTOMESON OSCILLATIONS

Consider neutral long-living LMs that interact with the SM leptons and quarks at the energies below the new physics scale Λ (e.g., the compositeness scale) dominantly through the effective four-fermion terms, unlike the ordinary sterile neutrinos with their Yukawa couplings to the leptons and the Higgs doublet. For the vector case with lepton number (L), lepton flavor and B conservation these four-fermion interactions at the first order in LM fields can be written as

$$\mathcal{L}_{\mathrm{int}} = \sum_{\psi_\ell, f, f'} \sum_{\alpha,\beta=L,R} \left[\frac{\epsilon^{\alpha\beta}_{ff'\psi_\ell}}{\Lambda^2} (\bar{f}_\alpha \gamma^\mu f'_\alpha)(\bar{\psi}_{\ell\beta}\gamma_\mu \ell^0_{M\beta}) + \frac{\tilde{\epsilon}^{\alpha\beta}_{ff'\psi_\ell}}{\Lambda^2} (\bar{\psi}_{\ell\alpha}\gamma^\mu f'_\alpha)(\bar{f}_\beta \gamma_\mu \ell^0_{M\beta}) \right] + \mathrm{H.c.}, \quad (1)$$

where ϵ and $\tilde{\epsilon}$ are the effective couplings (real couplings can work for the BG in this section), $\psi_\ell = \ell, \nu_\ell$ ($\ell = e, \mu, \tau$) is

[2]It has been shown in Refs. [38,39] that oscillations and decays of heavy sterile neutrinos are indeed two distinct sources for baryogenesis via leptogenesis, unlike some previous claims.

the SM lepton, f and f' denote either two quarks or two leptons (we take them from the same particle generation) such that the sum of the electric charges of f_α, $f'^\dagger_\alpha$ and $\psi_{\ell\beta}$ is zero, and ℓ^0_M is the neutral LM flavor state that is related to the mass eigenstates L^0_{Mi} by the mixing matrix U as

$$\ell^0_{M\alpha} = \sum_{i=1}^{n} U^\alpha_{\ell i} L^0_{Mi}. \tag{2}$$

LMs can be produced thermally from the primordial plasma. Once created ℓ^0_M oscillate and interact with ordinary matter. These processes do not violate the total lepton number $L^{\rm tot}$, defined as usual lepton number plus that of LMs. However, the oscillations of LMs violate CP and therefore their individual lepton numbers are not conserved.[3] Hence the initial state with all zero lepton numbers evolves into a state with $L^{\rm tot} = 0$ but nonzero individual lepton numbers of LMs.

At the temperatures below Λ scale LMs communicate their lepton asymmetry to the neutrinos and the charged leptons through the effective four-fermion interactions in Eq. (1). Suppose that the neutral LMs of at least one type come into thermal equilibrium before the time of EWSB $t_{\rm EW}$ at which sphalerons become ineffective, and those of at least one other type do not equilibrate by $t_{\rm EW}$. Hence the lepton number of the former (later) affects (has no effect on) the baryogenesis. As a result, the final baryon asymmetry after $t_{\rm EW}$ is nonzero. At the time $t \gg t_{\rm EW}$ all LMs decay into the leptons and the quarks (hadrons). For this reason they do not contribute to the dark matter in the Universe, and do not destroy the big bang nucleosynthesis.

The system of n types of singlet LMs with a given momentum $k(t) \propto T(t)$ that interact with the primordial plasma can be described by the $n \times n$ density matrix $\rho(t)$. In a simplified picture, this matrix satisfies the kinetic equation [40,48]

$$i\frac{d\rho}{dt} = [\hat{H}, \rho] - \frac{i}{2}\{\Gamma, \rho\} + \frac{i}{2}\{\Gamma^p, 1-\rho\}, \tag{3}$$

where Γ (Γ^p) is the destruction (production) rate, and the effective Hamiltonian can be written as

$$\hat{H} = V(t) + U\frac{\hat{M}^2}{2k(t)}U^\dagger, \tag{4}$$

where $\hat{M}^2 = {\rm diag}(M_1^2, \ldots, M_n^2)$ is the matrix of the squared LM mass eigenstates, and V is a real potential. (In the approximation of Boltzmann statistics, the last term in Eq. (3) is $i\Gamma^p$.) In general, evolution of LMs can be considered together with the evolution of the SM leptons using the methods of Refs. [41,49]. However, such precise numerical analysis is beyond the scope of the present consideration. In the following, we concentrate on the essentially different temperature dependence of the interaction rate for LMs and the sterile neutrinos, which makes the LM scenario more attractive for the experimentalists.

The cross sections for $2 \leftrightarrow 2$ reactions that contribute to the LM destruction rate can be written as

$$\sigma \equiv \sigma(a + b \leftrightarrow c + d) = \frac{C}{4\pi}\epsilon^2\frac{s}{\Lambda^4}, \tag{5}$$

where a, b, c and d denote the four interacting particles (f, f', ψ_ℓ and ℓ^0_M), $C = \mathcal{O}(1)$ is the constant that includes the color factor in the case of the interaction with quarks, s is the total energy of the process, and ϵ is the relevant coupling from Eq. (1). In the considered LM scenario, the cross section in Eq. (5) is proportional to s in contrast to the inverse proportionality in the case of BG from neutrino oscillations. The respective $2 \leftrightarrow 2$ scattering rate density at the high temperatures $M_i \ll T \ll \Lambda$ can be calculated as [23]

$$\begin{aligned}\gamma_s &= \frac{g_a g_b T}{32\pi^4}\int_0^\infty ds s^{3/2} K_1\left(\frac{\sqrt{s}}{T}\right)\sigma(s) \\ &= \frac{6C}{\pi^5} g_a g_b \epsilon^2 \frac{T^8}{\Lambda^4},\end{aligned} \tag{6}$$

where g_a is the number of internal degrees of freedom of the particle a, and K_1 is the Bessel function. Then the interaction rate that brings LMs into equilibrium (average destruction rate) can be written as

$$\Gamma \sim \epsilon^2 \frac{T^5}{\Lambda^4}. \tag{7}$$

The conditions that LMs of type L^0_i come into equilibrium before the time of the EWSB $t_{\rm EW}$, while LMs of type L^0_j do not, are

$$\Gamma_i(T_{\rm EW}) > H(T_{\rm EW}), \tag{8}$$

$$\Gamma_j(T_{\rm EW}) < H(T_{\rm EW}), \tag{9}$$

where the Hubble expansion rate H can be written as

$$H(T) \approx 1.66 g_*^{1/2}\frac{T^2}{M_{\rm Planck}}, \tag{10}$$

where $M_{\rm Planck} = 1.221 \times 10^{19}$ GeV is the Planck mass, and $g_* \sim 10^2$ is the number of relativistic degrees of freedom in the primordial plasma.

Remarkably, the rates in Eqs. (8) and (9) are suppressed by the fourth power of $T_{\rm EW}/\Lambda$ ratio with respect to the case of the BG via the sterile neutrino oscillations. For this

[3]For Majorana LMs the CP-violating scatterings can significantly effect the picture similarly to the case of sterile neutrinos [47].

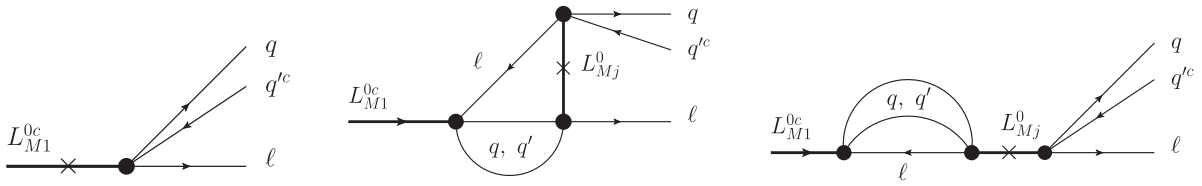

FIG. 2. Feynman diagrams for the discussed contributions to the CP asymmetry, where the $\times$ represents a Majorana mass insertion, the line direction shows either L or B flow, and the black bulbs represent subprocesses [a particular case of the leptoquark S_{0R} exchange is shown in Fig. 3 (left)].

reason, the couplings ϵ can be significantly larger than the Yukawa couplings of that sterile neutrinos. In particular, for $\Lambda \gtrsim 10$ TeV we have $\epsilon \gtrsim 10^{-4}$. Hence, the considered scenario of the BG via neutral LMs can be relevant for the LHC and the next collider experiments without unnatural hierarchy of couplings.

In the approximation of Eq. (3), the asymmetry transferred to the usual leptons by $t_{\rm EW}$ can be written as [40]

$$\frac{n_L - n_{\bar{L}}}{n_\gamma} = \frac{1}{2}\sum_j |S^M_j(t_{\rm EW}, 0)|^2_{CP-\rm odd}, \tag{11}$$

where the factor $1/2$ accounts for the photon helicities, and $S^M = U^\dagger S U$ is the evolution matrix in the mass eigenstate basis ($S(t, t_0)$ is the nonunitary evolution matrix corresponding to the operator $\hat{H} - (i/2)\Gamma$).

In the case of three LM mass states, the respective CP-violating effects should be proportional to the Jarlskog determinant [50] related to their mixing matrix U. However, extra LM mass states can enrich the picture of CP violation. Also additional CP-violating phases may come into play from the active neutrino sector (compare to Ref. [41]).

III. BARYOGENESIS FROM LEPTOMESON DECAY

Suppose that the neutral LMs are Majorana particles ($\ell^0_{MR} = \ell^{0c}_{MR}$). Then an analog of the usual LG can take place due to their out-of-equilibrium CP- and L-violating decays in the early Universe. The relevant terms among the B-, L- and lepton flavor-conserving LM interactions can be written as

$$\frac{\epsilon^{\alpha R}_{ff'\psi_\ell}}{\Lambda^2}(\bar{f}_\alpha \gamma^\mu f'_\alpha)(\bar{\psi}_{\ell R}\gamma_\mu \ell^0_{MR}) + \frac{\epsilon^{S}_{ff'\psi_\ell}}{\Lambda^2}(\bar{f}_R f'_L)(\bar{\psi}_{\ell L}\ell^0_{MR}) + \frac{\epsilon^{T}_{ff'\psi_\ell}}{\Lambda^2}(\bar{f}\sigma^{\mu\nu} f')(\bar{\psi}_{\ell L}\sigma_{\mu\nu}\ell^0_{MR}) + \text{H.c.}, \tag{12}$$

where the sum of the hypercharges of f, $f'^\dagger$ and ψ_ℓ is zero. To be more specific, in the following we consider the term

$$\frac{\lambda_{\ell i}}{\Lambda^2}(\bar{q}_\alpha \gamma^\mu q'_\alpha)(\bar{\ell}_R \gamma_\mu L^0_{Mi}), \tag{13}$$

where $\lambda_{\ell i} = \epsilon^{\alpha R}_{qq'\ell} U^R_{\ell i}$ is the complex parameter, and we used Eq. (2).

Consider the interference of the tree and two-loop diagrams[4] shown in Fig. 2, where L is violated by two units due to the Majorana mass insertion. The CP asymmetry that is produced in L^0_{M1} decays can be defined as

$$\varepsilon_1 = \frac{1}{\Gamma_1}\sum_\ell [\Gamma(L^0_{M1} \to \ell_R q_\alpha q'^c_\alpha) - \Gamma(L^0_{M1} \to \ell^c_R q^c_\alpha q'_\alpha)], \tag{14}$$

where the three-particle decay width is [53]

$$\Gamma_1 = \sum_\ell [\Gamma(L^0_{M1} \to \ell_R q_\alpha q'^c_\alpha) + \Gamma(L^0_{M1} \to \ell^c_R q^c_\alpha q'_\alpha)] \simeq \frac{1}{128\pi^3}(\lambda^\dagger\lambda)_{11}\frac{M_1^5}{\Lambda^4}, \tag{15}$$

with the mass M_1 of L^0_{M1}. For this CP asymmetry to be nonzero requires $\mathrm{Im}[(\lambda^\dagger\lambda)^2_{1j}] \neq 0$. Hence, at least two LM mass states are needed. In the case of quasidegenerate LM masses of $M_2 - M_1 \sim \Gamma_1/2 \ll M_1$, the self-energy graph gives the dominant contribution to the CP asymmetry that can be expressed in the same form [51,52] as in the usual resonant LG [33–35]. In the strong washout regime [54], the final $B - L$ asymmetry generated at $T \sim M_1$ is insensitive to any initial asymmetry at $T \gg M_1$. The respective condition for the decay parameter $K \equiv \Gamma_1/H(T = M_1) > 3$ translates into the limit of

$$(\lambda^\dagger\lambda)_{11} > 4 \times 10^{-7} \times \left(\frac{\Lambda}{10\ \text{TeV}}\right)^4 \times \left(\frac{1\ \text{TeV}}{M_1}\right)^3. \tag{16}$$

[4]Same two-loop self-energy graph was discussed in the resonant BG mechanisms of Refs. [51,52], where the baryon asymmetry is directly produced in the three-body decays of sterile neutrinos N. Although these mechanisms involve B-violating interactions of $QQQN$ type, they do not lead to fast proton decay due to the large values of N mass and the B-violating interaction scale of $\mathcal{O}(1)$ TeV. These mechanisms can be probed in the near future by the neutron-antineutron oscillations and other B-violating processes.

The discussed effective LM-quark-antiquark-lepton vertices can be economically realized, e.g., through exchange of the scalar $SU(2)_L$ singlet LQ S_{0R} with the weak hypercharge $Y = 1/3$ [55,56]. The relevant interaction terms in the Lagrangian can be written as

$$-\mathcal{L}_{\rm int} = (g_{ij}\bar{d}^c_R L^0_{Mi} + f_j \bar{u}^c_R \ell_R) S^j_{0R} + \text{H.c.} \qquad (17)$$

Then the above expressions are valid with the replacements $\lambda \to gf^*$ and $\Lambda \to M_{S_{0R}}$. In particular, typical values of the new couplings in Eq. (16), e.g., $|g| \sim |f| \sim 0.01$–0.1, can be interesting for the collider searches.

Notice that the new contributions to the CP asymmetry coming from the interferences among the tree and one-loop diagrams shown in Fig. 3 cancel each other unlike the more sophisticated case of Ref. [57] with the three types of interactions involved in the LG based on the three-body decays. However, the compositeness models with LQs, which have at least three types of interactions, can realize the LG of this kind from LM decays.

The final baryon asymmetry can be written as [58]

$$\begin{aligned}\eta_B &\equiv \frac{n_B - n_{\bar{B}}}{n_\gamma} = 7.04 \times \frac{n_B - n_{\bar{B}}}{s} \\ &= 7.04 \times \left(-\frac{28}{79}\right) \times \frac{n_L - n_{\bar{L}}}{s} \\ &= 7.04 \times \left(-\frac{28}{79}\right) \times \frac{\varepsilon_1 \kappa}{g_*}, \end{aligned} \qquad (18)$$

where n_B, n_L and n_γ are the baryon, lepton and photon number density, respectively, s is the entropy density, $\kappa \leq 1$ is the washout coefficient, and $-28/79$ is the sphaleron lepton-to-baryon factor. To exactly determine κ, one should solve the set of Boltzmann equations, which in the case of the resonant regime can be written as

$$\frac{dN_i}{dz} = -(D_i + S_i)(N_i - N_i^{\rm eq}), \qquad (19)$$

$$\frac{dN_{B-L}}{dz} = -\sum_i \varepsilon_i D_i (N_i - N_i^{\rm eq}) - N_{B-L} \sum_i W_i, \qquad (20)$$

where $z = M_1/T$ is a dimensionless variable, N_X $[N_X^{\rm eq}]$ (with $X = i$ and $B-L$) is the [equilibrium] number density of L^0_{Mi} and $B-L$, and the various reaction rates are denoted by the following factors: D_i for $L^0_{Mi} \to \ell q q'^c$ decays, S_i for the scatterings of $L^0_{Mi}\ell^c \to q q'^c$, $L^0_{Mi} q' \to \ell q$, etc., and W_i for the washout processes that include the scatterings and the inverse decays of $\ell q q'^c \to L^0_{Mi}$. It was shown in Ref. [51] for similar processes, which are generated by the operator $QQQN$ (instead of $L^0_M \bar{L} Q \bar{Q}$ operator in our model), that for the interesting parameter range of $M \sim 1$ and $\Lambda \sim 10$ TeV, the decay rate is much larger than the inverse decay rate and dominates over the scattering rates at $T \sim T_{\rm EW}$ as required for successful BG in the strong washout regime. Using the resonant CP asymmetry of

$$\varepsilon_i \sim \frac{\text{Im}\{[(\lambda^\dagger\lambda)_{ij}]^2\}}{(\lambda^\dagger\lambda)_{ii}(\lambda^\dagger\lambda)_{jj}} \frac{\Gamma_j}{M_j} \frac{M_i M_j}{M_i^2 - M_j^2} \sim \mu^{-1} \frac{\Gamma_1}{M_1}, \qquad (21)$$

the observed baryon asymmetry $\eta_B = (6.21 \pm 0.16) \times 10^{-10}$ [23] can be produced for the decay parameter of $K \sim 100$ and the degeneracy factor of

$$\mu \equiv \frac{M_2 - M_1}{M_1} \lesssim 10^{-6} \left(\frac{M_1}{1\ \text{TeV}}\right). \qquad (22)$$

Notice that the dependence of BG on the nonthermal production mechanism for the decaying neutral particles responsible for the BG was discussed in Ref. [52] for the model based on $QQQN$ operator.

IV. DISCUSSION AND CONCLUSION

A substantial feature of any successful BG scenario is consistency with the present bounds on the active neutrino masses. In the case of Majorana LMs among the discussed four-fermion interactions, the terms

$$\frac{\epsilon^S_{ff\nu_\ell}}{\Lambda^2}(\bar{f}_R f_L)(\bar{\nu}_{\ell L}\ell^0_{MR}) + \frac{\epsilon^T_{ff\nu_\ell}}{\Lambda^2}(\bar{f}\sigma^{\mu\nu} f)(\bar{\nu}_{\ell L}\sigma_{\mu\nu}\ell^0_{MR}) + \text{H.c.} \qquad (23)$$

can generate a two-loop contribution to the neutrino masses. For $f = q$, this contribution can be illustrated by the generic diagram in Fig. 4 (left), where the black bulbs represent some subprocesses. Its particular realization in a model with LQs is shown in Fig. 4 (right). The resulting neutrino mass can be estimated as

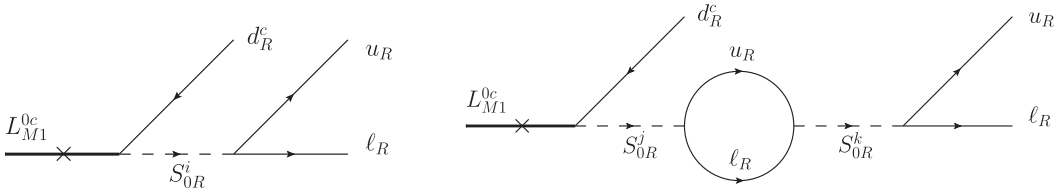

FIG. 3. Discussed Feynman diagrams in the model with scalar leptoquarks S^i_{0R}.

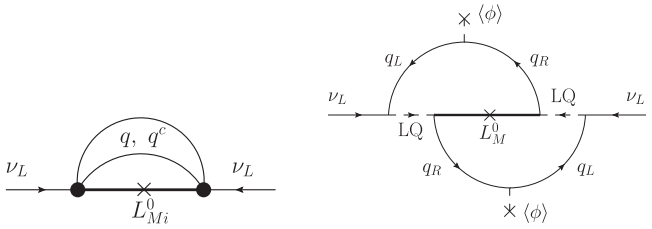

FIG. 4. Discussed diagrams for LM contribution to the neutrino masses.

$$m_{\nu_\ell} \sim \sum_i \frac{|\epsilon U_{\ell i}|^2}{(16\pi^2)^2} \frac{M_i^3 m_f^2}{\Lambda^4}, \tag{24}$$

where ϵ is a relevant coupling from Eq. (23). Then the present experimental upper bound on the neutrino mass of $m(\nu_e) \lesssim 2$ eV [59] can be easily satisfied for the discussed values of ϵ, M_i and Λ.

To conclude, we have introduced the two possible generic scenarios of low-temperature BG in the new class of models with LM states. The BG from LM decay can be realized if all LMs decay before the EWSB. In the case of relatively light and long-lived LMs, which do not all decay before the EWSB, the BG from LM oscillations may take place. One of the attractive features of this scenario is that the out-of-equilibrium condition is more relaxed with respect to the BG from the sterile neutrino oscillations. Namely, the constraint on the effective LM coupling ϵ is essentially weakened by the factor of $(\Lambda/T_{\rm EW})^2$ with respect to the strong constraints on the sterile neutrino Yukawas. For the contact interactions scale of $\Lambda \sim 10$ TeV, this factor is of $\mathcal{O}(10^3)$ and offers great prospects for the experimental searches of relevant LMs. Hence, an accurate examination of the allowed parameter spaces for the successful BG in the specific models with LMs is desirable in the next step.

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