B–L genesis by sliding inflaton

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Abstract. We propose a new mechanism of lepton (L) number asymmetry generation, hence offer an explanation of matter-antimatter imbalance when a significant amount of baryon number is later transformed from this L-number by known electroweak sphaleron mediated process. The basic theoretical framework is a recently proposed multiple scalar-tensor gravity that dynamically solves the cosmological constant problem. The L-asymmetry generation in one of two proposed scenarios is triggered by dynamical relaxation of scalar inflaton field towards the zero cosmological constant. CPT violation (C = charge conjugation, P = parity operation, T = time reversal) in the presence of a chemical potential gives the necessary time arrow, and lepton number violating scattering in cosmic thermal medium generates a net cosmological L-number via resonance formation. Another scenario is L-asymmetry generation from evaporating primordial black holes. These proposed mechanisms do not require CP violating phases in physics beyond the standard model: the new required physics is existence of heavy Majorana leptons of masses $10^{15} \sim 10^{17}$ GeV that realizes the seesaw mechanism. We identify the cosmological epoch of lepto-genesis in two scenarios, which may give the right amount of observed baryon to entropy ratio. It might even be possible to experimentally determine microscopic physics parameter, masses of three heavy Majorana leptons by observing astrophysical footprints of primordial black hole evaporation at specified hole masses.

Keywords: leptogenesis, modified gravity, dark energy theory, primordial black holes

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1 Introduction

Familiar baryo-genesis [1, 2], or lepto-genesis [3], scenarios are based on what may be called the delayed decay mechanism [4, 5].\textsuperscript{1} This mechanism uses short time span of the cosmic expansion during which heavy particles such as baryon number violating gauge and Higgs bosons, or right-handed Majorana leptons, decay without appreciable counter-balance due to inverse decays. The scenario typically requires the inequality relation between the thermally averaged decay rate $\Gamma$ and the Hubble rate $H$: $\Gamma > H > \Gamma_{\text{inv}}$ in order to block the inverse decay (with its thermally averaged rate $\Gamma_{\text{inv}}$). This inequality is usually regarded as the out-of-equilibrium condition. It was however later pointed out in [6, 7] that even balanced baryon or lepton number violating processes such as scattering among much lighter thermalized particles can generate the asymmetry provided that a finite chemical potential $\mu$ distinguishing particles from anti-particles is present. The chemical potential is introduced in grand canonical ensemble of statistical physics to cope with the case when the particle number is not exactly conserved. An advantage of this new scenario is that the asymmetry may in principle be accumulated during a longer time span. We shall use in the present work this mechanism within the framework of new cosmological models.

It was recently shown [8] that multiple scalar-tensor gravity solves major cosmological conundrums, at the same time answering the cosmological constant problem. In the proposed model the inflaton field has two stationary points; one at a spontaneously broken potential minimum around the Planck scale and the other at the field infinity where a properly defined dynamical cosmological constant vanishes. In the slow-roll to the potential minimum, inflation is realized to give a right amount of the observed spectral index. The field oscillation around this minimum then gives violent particle production towards thermalized hot big-bang [9–13]. The inflaton field oscillation simultaneously accompanies zero-point quantum fluctuation of another inflaton within the multiple scalar-tensor gravity, whose positive contribution may cancel the negative mass squared term in a wine-bottle type potential of

\textsuperscript{1}These are earliest reviews explaining some details of the delayed decay mechanism.
inflaton fields. Quantum fluctuation ultimately, but gradually, restores the broken symmetry of inflaton field system, when the inflaton field moves towards the field infinity of zero cosmological constant.

The long-time phase of dynamically relaxing inflaton is ideal for fostering generation of the $B-L$ asymmetry according to the idea of spontaneous $B-L$ genesis [6, 7], in particular, thermal genesis using a finite chemical potential that distinguishes particles from anti-particles. Evaporation of primordial black hole with the presence of a chemical potential is also a promising site of the asymmetry generation. We shall discuss these two possibilities and compare their results.

A particular mechanism we detail is based on a standard model extension adding heavy Majorana leptons (three $N_R$'s) responsible for the seesaw mechanism of small neutrino mass generation [14]. Lepton number violating scattering occurs in this model efficiently via resonance formation of $N_R$'s. Our scenario does not require CP violating phases in physics beyond the standard model. CPT violation is provided by a finite chemical potential which we calculate using standard model parameters in addition to the inflaton sliding velocity. Necessary physics inputs are all ready for implementation of this lepto-generation scenario. We shall be able to locate the energy scale of lepton number violation by taking the observed value of baryon to entropy ratio.

We use the natural unit of $\hbar = c = 1$ and the Boltzmann constant $k_B = 1$ throughout the present work unless otherwise stated.

2 Cosmological model realizing dynamical relaxation of cosmological constant; an overview

Since the preceding work [8] which we base on is not well publicized, we shall recapitulate main results relevant to our subsequent discussion of $B-L$ genesis in the present work. Derivation and detailed concepts that led to the preceding work are relegated to the original paper.

Towards construction of a unified cosmology the present author is strongly motivated by a belief that slow-roll inflation, late time acceleration (dark energy), and hopefully dark matter should be solved along with the cosmological constant problem [15]. Otherwise, one is not sure of at which cosmic epoch the cosmological constant should be fine-tuned; either the Planck epoch, or electroweak epoch, or QCD epoch, or some other epochs? It seems that the phenomenology of cosmology requires at least two nearly non-varying cosmological constants, one of the Planck order at inflation and the other of $(\text{meV})^2$ order at present: are they related or independent? There are two important concepts to realize the ultimate goal in our view: (1) dynamically varying cosmological constant, (2) restoration of spontaneously broken symmetry. We shall spell out these in some detail, adding two comments.

(1) An arbitrary, both in sign and magnitude, cosmological constant is introduced in gravity theories of multiple scalar $\chi$ fields different from general relativity, in which the Einstein-Hilbert action is changed to allow a lagrangian density of the form, $-F(\chi)G_N R$, with $R$ the Ricci scalar, $G_N$ the gravitational constant [16, 17]. The metric frame that has this form is called the Jordan frame. A positive definite function $F(\chi)$ is called the conformal function, which may be expanded into a polynomial of powers, $(\chi/M_P)^2$, $M_P = 1/\sqrt{16\pi G_N} \sim 2$.

Historically, the conformal coupling to Ricci scalar was introduced by these references, however without the potential term $V(\chi)$ and the cosmological constant.
$1.7 \times 10^{18}$ GeV. The inflaton field $\chi$ must have at least four real components. We assume both of conformal function and potential function $V(\chi)$ to possess a global O(4) symmetry.

Physical consequences of cosmology are best worked out in the Einstein metric frame that can be transformed from the Jordan frame by a Weyl rescaling, by which $F(\chi)$ factor is eliminated to unity in the Einstein frame, hence the non-varying gravitational constant emerges in this frame. An arbitrary cosmological constant $\Lambda$ introduced in the Jordan frame then becomes a dynamical variable of the form, $\Lambda/F^2(\chi)$, in the Einstein frame. This way the dynamical cosmological constant vanishes at the field infinity $\chi = \infty$. The standard model, or a grand unified model, is introduced as well in the Jordan frame, resulting in multiplication of some powers of $1/F(\chi)$ in the Einstein frame after the Weyl scaling. With time evolution of the inflaton $\chi$ fields, particle physics parameters apparently vary, which must be treated with great care, since the unit change of length and time is also present.

(2) Important epochs in cosmology can be regulated by a potential $V(\chi)$ for the inflaton field. The potential is assumed to have a wine-bottle type shape such that a spontaneous breaking occurs at a field value $\chi_m$. There are two corrections to this tree-level potential $V$: one is the Nambu-Goldstone (NG) kinetic repulsion [18], and the other is quantum fluctuation. NG kinetic contribution is incorporated as an effective potential term of centrifugal repulsion equal to constant $2/a(t)^6 \chi^2$, with $a(t)$ the cosmic scale factor, giving the shifting minimum position according to the cosmic expansion. The first important epoch occurs when the inflaton, starting from a field point $\chi > \chi_m$, settles down around the potential minimum at $\chi_m$. Due to the inflaton coupling to Higgs field $\lambda_H(|H|^2 - v^2)^2/(4F^2(\chi))$, $\lambda_H > 0$ (with $H$ the Higgs doublet, $v$ vacuum expectation value, and $\lambda_H$ dimensionless coupling) and other particle fields, we expect that a hot big-bang is realized using the mechanism of parametric amplification, as outlined in [9–13].

Effect of growing, though temporarily, quantum fluctuation, necessarily accompanying particle production, may drastically change the nature of O(4) symmetry among four real inflaton fields. The original negative mass squared term $\propto -\chi^2 \chi^2$ in $V(\chi)$ receives a positive-definite correction $\propto \langle \phi^2 \rangle$ where $\phi$ is an orthogonal component to the dominant $\chi$ inflaton field. This gives rise to O(4) symmetry restoration, lowering a wall height right to $\chi_m$ and finally making the field infinity the global potential minimum. The inflaton $\chi$ field then starts to roll down to the potential bottom at which dynamical cosmological constant vanishes. It is important that the field zero $\chi = 0$ can never be reached due to that the NG repulsion makes the field origin an infinitely high wall.

(3) Caveat. The importance of the metric frame difference in physical interpretation seems to have been frequently overlooked in the past literature. Many popular dark energy models introduce non-minimal kinetic scalar terms of the form $G(\chi)(\partial \chi)^2$ in the Einstein frame. These models may be regarded to belong to the same class of models as the model of [8], when one rescales back to the Jordan frame according to Weyl. But quite often it is not mentioned in the literature in which frame the standard particle physics lagrangian and a cosmological constant are introduced. The cosmological constant problem becomes a dynamical issue, only when these are introduced in the Jordan frame, and not in the Einstein frame. The Jordan frame seems more fundamental as a starting point if four dimensional field theory descends from higher dimensional Kaluza-Klein unification or superstring theories.

(4) Example of compared metric frame difference. The metric frame differences have been analyzed in great detail for simple models in [19]. In the Jordan-Brans-Dicke theory [16, 17]
added by a cosmological constant $\Lambda$, attractor solutions that give cosmic scale factor $\propto t^{1/2}$
dynamical cosmological constant $\propto \Lambda/t^2$ are found in the Einstein frame. This model
is equivalent to a quintessence model [20–22] of potential $\propto \chi^{-2}$ in the Einstein frame, or
$F(\chi) \propto \chi^2$ in the Jordan frame. Surprisingly, this solution corresponds to attractor of the static Minkowski spacetime in
the Jordan frame. All attractor solutions in the Einstein frame are shown [19] to exhibit the dynamical relaxation of cosmological constant.

We need to give some more details of our model for subsequent discussion of leptogenesis. The potential and conformal functions
are both quartic in the Jordan frame,

\begin{equation}
F(\chi) = 1 + \xi_2 \left( \frac{\chi}{M_P} \right)^2 + \xi_4 \left( \frac{\chi}{M_P} \right)^4, \quad \xi_i > 0, \tag{2.1}
\end{equation}

\begin{equation}
V^{(J)}(\chi) = \frac{g}{4} \left( \chi^2 - \frac{m^2}{g} \right)^2 + 2M_P^2\Lambda + (\text{NG centrifugal repulsion terms}), \quad g > 0. \tag{2.2}
\end{equation}

This is the simplest scheme for successful realization of ideas stated above. The basic inflaton
field equation (when a single-field description is applicable) is given by

\begin{equation}
\ddot{\chi} + 3\frac{\dot{a}}{a}\dot{\chi} - \frac{\partial \chi}{F}F \dot{\chi}^2 = -\partial \chi V^{(E)}_{\text{eff}}(\chi), \tag{2.3}
\end{equation}
suppressing contributions from spatially inhomogeneous inflaton fields. We refer to the original paper [8] for the precise definition of how the effective potential $V^{(E)}_{\text{eff}}(\chi)$ is related to $V^{(J)}(\chi)$ and $F(\chi)$, but its approximate form relevant to our subsequent discussion is given shortly.

The slow-roll phase of inflation occurs in the field region that general relativity approximately holds, namely at $F(\chi) = 5$. General relativistic analysis of slow-roll condition [23] and calculation of the spectral index along with the tensor to scalar mode ratio [24] can be applied to constrain the potential derivative and curvature. We confirmed [8] that the slow-roll inflation is realized consistently with observations [25].

To the rest of discussion on B–L genesis, it is important to give result of time evolution
of inflaton field in radiation dominant (RD) epoch. During RD epoch the inflaton field essentially
follows a single-field equation, and the effective potential needed for a large $\chi$ field is

\begin{equation}
V_{\text{eff}}^{(E)}(\chi) \approx \frac{M_P^4}{5} \left( \text{constant} - \frac{g}{\xi_4} \ln \frac{\chi}{\chi_*} \right), \tag{2.4}
\end{equation}

where $\chi_*$ is a starting point of field roll-down to infinity after the symmetry restoration, and
its precise value is irrelevant to our discussion.

The field equation in RD epoch is approximately

\begin{equation}
\ddot{\chi} + \frac{3}{2t}\dot{\chi} = \frac{g}{5\xi_4} \frac{M_P^4}{\chi}. \tag{2.5}
\end{equation}

One arrives at the following solution, supported by some limited numerical simulations,

\begin{equation}
\chi(t) = \sqrt{\frac{2g}{15\xi_4}} M_P^2 t. \tag{2.6}
\end{equation}

\footnote{This is a standard textbook of modern cosmology. Section 4.2 explains how the condition for the slow-roll inflation is derived in single-field inflation models based on general relativity.}
Namely, the first term $\dot{\chi}$ and $\partial \chi F/F$ term are negligible compared to the Hubble term of right-hand side of eq. (2.3). Since the redshift factor is defined by $z + 1 = a(t_0)/a(t)$ with $t_0$ the present cosmic time, this gives a simple relation, $(H_0 = \text{the present Hubble rate}),$

$$\chi(t) = \chi(t_0)(z + 1)^{-2}, \quad \chi(t_0) \approx \sqrt{\frac{T}{30}} \sqrt{\frac{\varphi}{\xi_4}} \frac{M_P^2}{H_0} \sim 0.31 \times 10^{60} \frac{\varphi}{\xi_4} M_P.$$  \hspace{1cm} (2.7)

Thus, in RD epoch the inflaton field $\chi$ increases like $\propto T^{-2}$ with the cosmic temperature, as the cosmic time increases. At later times after recombination the field increases more slowly, typically $\propto (t/H_0)^2$, but this behavior is irrelevant to our discussion of lepto-genesis. The $F-$factor in RD epoch at the field infinity is approximately

$$F^{-2} \approx \left( \frac{\xi_4}{M_P} \right)^4 \sim (z + 1)^{16}.$$  \hspace{1cm} (2.8)

It has been argued that cold dark matter is generated as spatially inhomogeneous components of inflaton fields generated from thermal cosmic medium [8]. The ratio of gravitational strength to a clump mass $M$ made of $\chi$ cold dark matter, $G_N M^2$, increases tremendously, as discussed in [26], giving a rational of considering the subsequent scenario of black hole evaporation for the asymmetry generation. A typical growth of $G_N M^2$ is given by $\propto (z + 1)^{16}$. One would expect that these strongly bound systems merge into heavier black holes, emitting primordial gravitational waves. On the other hand, if the clump is made of ordinary nucleons, the ratio grows $\propto (z + 1)^{16 \epsilon}$ where $\epsilon$ was introduced for consistency with nucleo-synthesis and requirement of local gauge invariant field theories [26]. Even if there exist plenty of primordial black holes, it is not clear how much fraction of these is made of $\chi$ matters. In our black hole evaporation scenario we shall have to assume that there exist sufficiently large amount of $\chi$ primordial black holes.

3 Generation of chemical potential for particle anti-particle asymmetry

Inflaton field $\chi$ may couple with the B−L number current, to give an effective lagrangian [6, 7] of the form,

$$\mathcal{L}_{B−L} = \eta \partial \mu \chi J^\mu_{B−L} = -\eta \chi \partial \mu J^\mu_{B−L} + \text{total derivative},$$  \hspace{1cm} (3.1)

where $J^\mu_{B−L}$ is the B−L number four-current. $\eta$ has a negative mass dimension of $-1$, and is suppressed by a new physics scale beyond the standard particle physics model. When the B−L number is violated with $\partial \mu J^\mu_{B−L} \neq 0$, this CP-odd operator may give rise to the B−L asymmetry. In a spatially homogeneous FRW (Friedmann-Robertson-Walker) metric, $ds^2 = dt^2 - a^2(t)(d\vec{x})^2$, this effective term is equivalent to

$$\mathcal{L}_{B−L} = \mu_{B−L} n_{B−L}, \quad \mu_{B−L} = \eta \dot{\chi}.$$  \hspace{1cm} (3.2)

The quantity $\mu_{B−L}$ is the chemical potential that distinguishes particles of positive B−L from anti-particles of negative value. The asymmetric number density, namely difference between particles and anti-particles, is given by

$$\frac{1}{3} \mu_{B−L} T^2,$$  \hspace{1cm} (3.3)
Figure 1. Effective $\chi \bar{\ell} (\ell = e, \mu, \tau)$ vertex diagram for chemical potential $\mu_{B-L}$. Wavy circulating line is neutral Higgs boson $H$.

for the chemical potential $\mu_{B-L} \ll T$. In the model of [8] the inflaton velocity $\dot{\chi}$ is maximally large and of the Planck order, $O(M_P^2)$, as seen from (2.6). The generated $B-L$ number at early epochs of cosmological evolution is re-shuffled via sphaleron [27] mediated process that later occurs [28, 29].

It was pointed out that if the generation of chemical potential occurs in the phase of scalar field oscillation around a potential minimum, the generated asymmetry is wiped out [30]. But our proposed mechanism here occurs in the phase of monotonic inflaton sliding of $\dot{\chi} > 0$, and this criticism does not apply. The monotonic roll down to the field infinity is an essential element of resolving the cosmological constant problem in the scheme of [8].

The chemical potential in statistical physics is defined in the hamiltonian formalism. In [31] it was shown that the proper chemical potential thus defined may differ from $\dot{\theta}$ used here and in the literature [6, 7]. But the use of $\dot{\theta}$ for discussion of $B-L$ genesis remains valid, hence we shall keep this terminology for a mere convenience.

This mechanism of $B-L$ genesis has features distinct from the familiar delayed decay mechanism [1, 2]. In general this makes the generated asymmetry larger because one does not need CP violating phases that necessarily arises in higher orders of perturbation [4]. The need for time arrow is provided by a non-vanishing time derivative of inflaton sliding velocity $\mu_{B-L} \propto \dot{\chi} \neq 0$. The time span or temperature span is given by the temperature range of inequality $T > T_D$ when this mechanism works, and it is only limited by the decoupling temperature $T_D$.

The chemical potential is calculated from a triangular one-loop diagram, as shown in figure 1, for the effective vertex of three external fields, $\chi \bar{\ell} l$ with $l$ one of charged leptons, $e, \mu, \tau$. Its probability amplitude has the form in the low energy limit,

$$A_0 = \frac{i}{2(4\pi)^2} c_{\chi H} y_H^2 m_H \bar{\ell} \int_0^1 dx \frac{(1-x)^2}{m_H^2 (1-x) + m_l^2 x},$$

(3.4)

where $m_H$ is the Higgs boson mass, $m_l$ is a lepton mass, $y_H = m_l/v, v = 254\text{GeV} \times \sqrt{2}$ is the Yukawa coupling. The parameter $c_{\chi H}$ arises from inflaton coupling to Higgs boson of a vertex form $\chi HH$, [8]

$$V_{\chi H}^{(E)} = \frac{\lambda_H}{4 F^2(\chi)} \left(|H|^2 - v^2\right)^2, \quad c_{\chi H} = 8\lambda_H \frac{v^2}{\chi} \left(\xi_4 \frac{\chi^4}{M_P^4}\right)^{-2}. \quad (3.5)$$

The quantity $\chi/M_P$ evolves with cosmic time $t$ as given by (2.6); $\chi/M_P = O(M_P t)$ in RD epoch. Parameters $g, \xi_4$ are dimensionless coupling constants of order unity given in (2.1)
and (2.2). The time evolution of inflaton field is given by the redshift dependence \( \propto (z + 1)^{-2} \) of sliding inflaton \( \chi(t) \). The redshift is defined by the ratio of scale factor to its present value, \( z + 1 = a(t_0)/a(t) \) and is proportional to the cosmic temperature ratio \( T/T_0 \). The integral is simplified for a mass relation \( m_H \gg m_l \):

\[
\int_0^1 dx \frac{(1 - x)^2}{m_H(1 - x) + m_l^2} \sim \frac{1}{2m_H^2}.
\]

The time evolution of the redshift dependence is given by the redshift dependence \( \propto (z + 1)^{-2} \) of sliding inflaton \( \chi(t) \). The redshift is defined by the ratio of scale factor to its present value, \( z + 1 = a(t_0)/a(t) \) and is proportional to the cosmic temperature ratio \( T/T_0 \). The integral is simplified for a mass relation \( m_H \gg m_l \):

\[
\int_0^1 dx \frac{(1 - x)^2}{m_H(1 - x) + m_l^2} \sim \frac{1}{2m_H^2}.
\]

The probability amplitude is then given by

\[
A_0 \sim \frac{i}{4(4\pi)^2} \frac{c_{\chi H}}{m_H^2} \sum y_{HI}^2 m_l \bar{l}.
\]

It is interesting that this formula depends on known parameters of electroweak theory except \( c_{\chi H} \) of (3.5). The dominant contribution is from a largest \( y_{HI} \), hence from the coupling to the heaviest charged lepton, \( \tau \) pair. The chemical potential is thus

\[
\mu_L = \frac{c_{\chi H}}{4(4\pi)^2} \frac{m_\tau^2}{m_H^2} v^2 \dot{\chi}, \quad \frac{c_{\chi H}}{v^2} \dot{\chi} = 8\lambda_H(\xi_4)^{-2} \left( \frac{\chi}{M_P} \right)^{-9} \sqrt{\frac{2g}{15\xi_4}} M_P.
\]

We replaced the general notation \( \mu_{B-L} \) in previous formulas by the leptonic chemical potential \( \mu_L \) to better fit the present model. This chemical potential is extremely sensitive to the inflaton rolling of \( \chi \), hence is rapidly decreasing function of redshift factor \( \propto (z + 1)^{-18} \propto T^{18} \). The high power of this behavior arises from \( c_{\chi H} \) of (3.5), and the inflaton linear velocity \( \dot{\chi} \) in \( \mu_L \) gives no redshift dependence as clear from (2.6).

### 4 Lepton number violating scattering via resonance

We shall discuss in the following section two possible scenarios of lepton number generation: firstly thermal genesis and secondly primordial black hole evaporation. We need a common source of lepton number violating process in these two scenarios, and for this matter we need physics beyond the standard particle physics theory.

We consider SO(10) related models [14, 32]. More specifically, it is not even necessary to completely specify SO(10) models, and any SO(10) model that introduces the Majorana \( N_R \) mass term violating the lepton number and Dirac-type mixing arising from its Higgs coupling to left-handed neutrino field of the form,

\[
\left( N_R^T \mathcal{M} i\sigma_2 \sigma \cdot \partial N_R + \text{(h.c.)} \right) + \left( N_L^T m\nu_L + \text{(h.c.)} \right),
\]

is all we need. We used two-component spinor fields with \( \sigma = (1, \vec{\sigma}) \). Diagonal elements \( M_N \) of the \( 3 \times 3 \) matrix \( \mathcal{M} \) are assumed to be much larger than the standard model scale of 1 TeV, presumably close to grand unification scale or the Planck scale \( M_P \). For a single flavor case the mass diagonalization of the neutral lepton sector provides the seesaw mechanism [14], giving a tiny Majorana mass of order, \( m_\nu = m^2/M_N \), for ordinary neutrinos. Heavy Majorana particles stay with their masses of order \( M_N \). We thus assume only a minimum extension of standard model to accommodate B–L violation based on already experimentally established neutrino oscillation.

Actually, existing neutrino oscillation data [33] suggest that there may be a hierarchy of three neutrino masses, the heaviest having a larger mass than of order 50 meV, the next
heaviest of order $10\text{meV}$, and the lightest mass being unknown. There is a plenty of room accommodating the situation like this one, since many parameters, three diagonal elements $M_N$, $m$’s, and the mixing matrix elements, combine to define ordinary neutrino masses.

A finite chemical potential supported by lepton number violating process gives a basis of asymmetry generation. As the major process of lepton number violation that occurs till later epochs one can think of inelastic scattering of charge exchange, $l^+ W^+ (H^+) \rightarrow \bar{l}^+ W^- (H^-)$, with $W^\pm (H^\pm)$ charged weak gauge (Higgs) boson, with its Feynman diagram shown in figure 2. The charge exchange scattering amplitude is proportional to the Majorana mass $M_N$, while the associated elastic scattering amplitude is proportional to kinetic term. Prior to the electroweak phase transition the longitudinal part of $W^\pm$, namely the absorbed part $H^\pm$ in the minimum Higgs doublet model, contributes. Since the cross sections for lepton number decreasing $l \rightarrow \bar{l}$ and increasing $\bar{l} \rightarrow l$ processes are equal at the tree diagram level, the net generated asymmetry is given by the difference of thermal phase space factors with different signs of the chemical potential, as given by (3.3).

The probability amplitude of lepton-number violating process $l^+ W^+ (H^+) \rightarrow \bar{l}^+ W^- (H^-)$ via Majorana $N_R$ exchange is, in the free space,

$$A_L = -i g_{HN}^2 M_N \psi^c \left( \frac{1}{s - M_N^2} + \frac{1}{u - M_N^2} \right) \psi_u,$$

in terms of invariant squared momentum sum, $s = (p_l + p_H)^2$, $u = (p_l - p_{\bar{l}})^2$. $\psi_u$ and $\psi_v$ are four-spinor plane wave functions of lepton and anti-lepton. This amplitude is for a single $N_R$, and one must sum over three species of Majorana leptons. $g_{HN}$ is the bilinear Higgs coupling to $\bar{\nu}_L N_R$. From this probability amplitude one calculates the cross section in the free space, and takes the thermal average over initial leptons and Higgs bosons at temperature $T$.

Thermally averaged cross sections differ depending on whether the temperature is in the mass region, $T \approx M_N$, or outside this region. The thermal average is calculated and numerically illustrated in appendix. We shall mention main results here. The low energy limit $T/M_N \rightarrow 0$ has a finite value independent of temperature, of order $g_{HN}^2 M_N^2$ as given by (A.8), while the high energy limit $T/M_N \rightarrow \infty$ decreases as a temperature power $\propto 1/T^4$ as given by (A.9). The most important is in the resonance energy region at $s = M_N^2$. The resonance temperature is estimated as $T_R \approx M_N/3$, as given by (A.6). The resonance width is $\Delta T \approx 0.086 \gamma_N$ in terms of natural + thermal width $\gamma_N$; a quantity not easy to calculate at present. The cross section in the narrow width limit behaves as a Lorentzian function almost like a Dirac delta-function, giving a large resonance cross section $\propto 1/\Delta T$. The behavior of thermally averaged cross section is illustrated for large widths (numerically difficult to work out narrow width cases) in figure 3.
5 Thermal generation and primordial black hole evaporation

5.1 Thermally generated lepton asymmetry

Before we discuss the main issue of this section, we first note that the inflaton time evolution in section 2 is conveniently written in terms of the Planck energy $M_P \sim 1.7 \times 10^{27}$ eV, the present Hubble rate $H_0 \sim 1.6 \times 10^{-33}$ eV, the present microwave temperature $T_0 \sim 2.35 \times 10^{-4}$ eV, and the temperature in radiation dominated (RD) epoch $T$:

$$\left( \frac{\chi}{M_P} \right)^{-1} = \sqrt{\frac{30}{\lambda_4 T}} \frac{H_0 T^2}{g M_P T_0^3} \sim 9.5 \sqrt{\frac{\xi_4}{g}} \left( \frac{T}{10^{17} \text{GeV}} \right)^2. \quad (5.1)$$

First, we discuss the temperature region in which lepton number violating scattering works. In cosmic thermal medium $L$-violating scattering process decouples at a temperature $T_D$ derived from the relation $\Gamma_L = 3H$ with $H \propto T^2/M_P$ the Hubble rate. The high power dependence of $\Gamma_L \propto \mu_L \propto (\chi/M_P)^{-9}$ makes the decoupling temperature depend on parameters of the model unusually,

$$T_D = 1.4 \times 10^{16} \text{GeV} \left( \frac{\gamma N}{M N} \right)^{1/16} \left( \frac{\xi_4^2}{\lambda H} \sqrt{\xi_4} \right)^{1/16} N_{\text{eff}}^{1/32}. \quad (5.2)$$

We have used the resonance formation cross section $\sigma$ in the rate formula $\Gamma_L = \sigma \mu_L T^2/6$, calculated in appendix; eqs. (A.5) and (A.10). $\gamma N$ is the resonance width. The right-hand side value is insensitive to details of model parameters. There are three resonances, $\sim M_N^i/3$, $i = 1, 2, 3$ corresponding to three diagonal mass matrix elements. The decoupling temperature $T_D$ should be set below the smallest resonance temperature $T_R \approx M_N/3$ in order to fully exploit large resonance contributions.

We next turn to the amount of generated asymmetry. The instantaneous $L$-asymmetry generation rate $\Gamma_L$ is given in terms of thermally averaged cross section $\sigma(T)$,

$$\Gamma_L(T) = \sigma(T) \frac{\mu_L T^2}{6}. \quad (5.3)$$

As already stated, this rate may accumulate during a long time span, and the accumulated asymmetry is given by

$$\mathcal{Y} = \int_0^{T_D} dt \Gamma_L(T) = 2 \int_{T_D}^{\infty} dT \frac{t(T)}{T} \Gamma_L(T) = \frac{2}{\pi} \sqrt{\frac{45}{N_{\text{eff}}}} M_P \int_{T_D}^{\infty} dT \frac{\Gamma_L(T)}{T^3}. \quad (5.4)$$

Here $N_{\text{eff}}$ of order 100 is the effective massless degrees of freedom contributing to cosmic energy density. This formula is valid in RD epoch in which $t \propto M_P / T^2$.

Assuming dominant contribution to arise from the resonance region $T = T_R \approx M_N$, one may approximate by the integral here,

$$\int_{T_D}^{\infty} dT \frac{\Gamma_L(T)}{T^3} \approx \frac{1}{6} \sigma_R \Delta T \left( \frac{\mu_L}{T} \right)_{T = T_R}. \quad (5.5)$$

From formulas in the preceding sections, we may write

$$\sigma_R \Delta T \approx \frac{8.77}{8\pi^4} \frac{y_{HN}^4 M_N}{\gamma N T_R} \sim 0.84 \times 10^{-2} \frac{y_{HN}^4}{\gamma N} \quad (5.6)$$

$$\left( \frac{\mu_L}{T} \right)_{T = T_R} \approx 3.2 \times 10^{-9} \frac{1}{\sqrt{N_{\text{eff}}}} \sqrt{\frac{g}{\xi_4}} \left( \frac{10^{16} \text{GeV}}{M_N} \right)^2. \quad (5.7)$$
From these we derive

\[ Y \approx 3.3 \times 10^{-9} \frac{y_{HN}}{N_{\text{eff}}} \left( \frac{10^{16} \text{GeV}}{M_N} \right)^\frac{3}{2} \frac{g}{\xi_4 \sqrt{\lambda_H}}. \]  

(5.8)

Due to three resonances this number is essentially tripled. The lepton/entropy ratio is down by \( n_L/s \) with the entropy density \( s = 2(s_B + 7s_F/8)T^3/3 \). With the electroweak baryon non-conservation \cite{28} taken into account, the remaining baryon to entropy ratio is estimated to be \( \mu_L/T_D \) times \( O(0.1) \) \cite{29} \cite{34}. There is thus a possibility to explain the observed baryon asymmetry. Our result \( Y \) weakly depends on the Majorana mass, \( \propto M^{-3}_N \).

Resonance formation kinematically requires high ambient temperatures larger than the mass of heavy Majorana particle; \( T > O(M_N) \). On the other hand, lower reheating temperatures are demanded by the density perturbation generated at inflation. The decoupling temperature given by (20) assuming resonance formation appears a bit large from this point, although a more precise formula of inflaton time evolution than given by (19) is to be used to arrive at a definite conclusion. In view of this we explored the possibility of low ambient temperatures taking lepton number violating scattering off the resonance region. Using low temperature limit of scattering cross section, one derives decoupling temperature,

\[ T_D = 0.95 \times 10^{15} \text{GeV} N_{\text{eff}}^{1/36} y_{HN}^{-2/9} \left( \frac{g^2}{\xi_4^2 \lambda_H} \right)^{1/9}. \]  

(5.9)

This presumably clears the condition of low reheat temperature for the density perturbation. The generated baryon asymmetry is calculated as

\[ Y \approx 3.8 \times 10^{-15} N_{\text{eff}}^{-1/2} \left( \frac{10^{16} \text{GeV}}{M_N} \right)^2 \left( \frac{T_D}{10^{16} \text{GeV}} \right)^{18} \frac{\lambda_H \xi_4^2}{y^4}. \]  

(5.10)

This magnitude is too small compared to the observed value. An intermediate temperature region between the resonance formation and the low temperature limit may give a good solution, but we postpone this study in view of missing improved treatment of inflaton time evolution.

5.2 Asymmetry generated by primordial black hole evaporation

We turn to the scenario of black hole evaporation. The idea of using primordial black hole (PBH) evaporation to generate the baryon asymmetry is old \cite{36} \sim \cite{41}. We mention a further benefit of our approach over the old scenario: the inflaton field in the model of \cite{8} may gravitationally collapse to readily form primordial black holes due to stronger gravity \cite{26} than general relativity would predict. This way primordial black holes can make up the majority of cold dark matter, while some fraction may generate the lepton asymmetry.

We first recapitulate main features of black hole evaporation according to \cite{35–37}.

Black holes evaporate thermally with a temperature known as Gibbons-Hawking temperature related to the hole mass \( M_{\text{BH}} \),

\[ T_{\text{BH}} = \frac{16\pi M_P^2}{M_{\text{BH}}} \sim 2.65 \times 10^{17} \text{GeV} \frac{mg}{M_{\text{BH}}}. \]  

(5.11)

\footnotetext{6}{The original idea of baryo-genesis using primordial black hole evaporation goes back to \cite{39} and \cite{41}.}
It is a well-publicized remarkable feature that smaller black holes evaporate heavier particles more than lighter ones, but at a temperature very close to the lifetime it is not entirely clear that this description remains true due to many fundamental issues related to information loss paradox. Due to this Hawking radiation, the hole mass changes with time according to

\[ M_{BH}(t) = M_{BH} \left(1 - \frac{t}{\tau_{BH}}\right)^{1/3}, \quad \tau_{BH} = \frac{20}{\pi N_{\text{eff}} M_{P}^{3}} \approx 27 \gamma \left(\frac{M_{BH}}{10^{12} \text{ gr}(N_{\text{eff}}/100)^{1/3}}\right)^{3}, \quad (5.12) \]

with \( M_{BH} \) the initial hole mass at \( t = 0 \). Although this is not the exponential decay law, one often calls \( \tau_{BH} \) the black hole lifetime. The mass of PBH’s evaporating now is estimated by equating \( \tau_{BH} \) to the present age \( \sim 13.5 \text{ Gy} \), to give \( M_{BH} \sim 8.0 \times 10^{14} \text{ gr}(N_{\text{eff}}/100)^{1/3} \).

Observation of \( \sim 1 \text{ year} \) gives a history of PBH evaporation for a hole mass \( \sim 3 \times 10^{11} \text{ gr} \).

There is an important kinetic constraint for successful PBH evaporation scenario: the evaporation rate of order \( 1/\tau_{BH} \) must be larger than the cosmic expansion rate given by the Hubble rate. Equivalently this constraint is given by \( \tau_{BH} \ll t \) (cosmic time). This leads to the constraint on cosmic temperature \( T \) in RD epoch,

\[ \frac{20 M_{BH}^{3}}{\pi N_{\text{eff}} M_{P}^{4}} > \sqrt{\frac{45 M_{P}}{\pi^{2} N_{\text{eff}} T^{2}}}. \quad (5.13) \]

This gives a maximum cosmic (not black hole) temperature

\[ T < 3.15 \times 10^{18} \text{ GeV} \left(\frac{M_{P}}{M_{BH}}\right)^{3/2} \left(\frac{N_{\text{eff}}}{100}\right)^{1/4}. \quad (5.14) \]

For 1 mg mass of the hole, the right-hand value is \( 1.7 \times 10^{17} \text{ GeV}(N_{\text{eff}}/100)^{1/4} \). It is interesting that this temperature range covers the energy range of resonance production discussed in the preceding subsection.

The asymmetry instantaneously generated by a single black hole is given by ratio of asymmetric number density to entropy density:

\[ \frac{c_{\mu L}}{\tau_{BH}} = c \mu_{L} G_{N} M_{BH}, \quad c \approx \frac{1}{2 s_{B} + 7 s_{F}/4} = O(0.01). \quad (5.15) \]

Although the rate of L-number violating scattering \( \propto \sigma \) of the preceding section does not enter in this formula, its relevance becomes evident when one discusses at which PBH temperature \( T_{BH} \) the asymmetry generation occurs. The baryon to photon number ratio in the universe is this number times the fraction \( f \) of PBH abundance.

The instantaneous L-asymmetry generation rate accumulates during the whole epoch of PBH evaporation, and this gives the accumulated averaged asymmetry,

\[ \xi_{BH} = \frac{4 c}{3} G_{N} M_{BH} \int_{0}^{t_{BH}} dt \frac{1}{\tau_{BH}^{1/3}} \mu_{L} \left(1 - \frac{t}{\tau_{BH}}\right)^{1/3}. \quad (5.16) \]

The chemical potential given in (3.8) contains the parameter combination \( c_{\mu H} \xi_{4}^{-2} \) dependent on relevant temperature of medium environment. This can be expressed using (5.1),

\[ \frac{c_{\mu H} \xi_{4}}{v^{2} M_{P}} = 8 \cdot 9.5^{9} \lambda_{H} \xi_{4}^{-2} \sqrt{\frac{2 g}{15 \xi_{4}}} \left(\frac{\xi_{4}}{g}\right)^{9/2} \left(\frac{T_{BH}}{10^{17} \text{ GeV}}\right)^{18} \sim 1.8 \times 10^{9} \lambda_{H} \xi_{4}^{-2} \left(\frac{\xi_{4}}{g}\right)^{4} \left(\frac{T_{BH}}{10^{17} \text{ GeV}}\right)^{18}, \quad (5.17) \]

\[ \mu_{L} = 5.9 \times 10^{5} M_{P} \left(\frac{\xi_{4}}{g}\right)^{4} \left(\frac{T_{BH}}{10^{17} \text{ GeV}}\right)^{18} \sim 1.3 \times 10^{35} \text{ GeV} \left(\frac{\xi_{4}}{g}\right)^{4} \left(\frac{\text{mg}}{M_{BH}}\right)^{18}, \quad (5.18) \]
using \((m_\tau/m_H)^2 \sim 2.0 \times 10^{-4}\) and (5.11). This gives the accumulated L-asymmetry of order,

\[ \mathcal{Y}_{BH} = cG_N M_{BH}/\mu L \sim 3.3 \times 10^{-10} \frac{\xi_4}{16\pi} \left( \frac{50 \text{ mg}}{M_{BH}} \right)^{17}. \]  

(5.19)

Due to the high power \(\propto M_{BH}^{-17}\), there is a strong preference of evaporating PBH mass. There is a narrow PBH mass region around a few to several tens of mg that gives a right amount of baryon to entropy ratio of order \(10^{-10}\). Generated lepton asymmetry from PBH evaporation prior to electroweak epoch is later converted to baryon asymmetry, although detection of the remaining lepton asymmetry is very difficult.

A problem inherent to the asymmetry generation based on primordial black hole evaporation is that it is difficult to reliably predict the mass distribution of potential candidates of primordial black holes. It is possible for a portion of primordial black holes to evaporate and the rest to remain as dark matter, but how much is left as dark matter is difficult to estimate.

We point out an exciting possibility. Suppose that the \(N_R\) resonance temperature \(T_R\) coincides with the evaporating black hole temperature \(T_{BH}\), which gives a relation,

\[ M_N = 1.6 \times 10^{16} \text{ GeV} \frac{50 \text{ mg}}{M_{BH}}. \]  

(5.20)

The preferential selection of PBH mass thus favors a particular \(N_R\) mass, thus an important microscopic parameter may be determined, though otherwise extremely difficult to measure, if one confirms experimentally PBH evaporation in this mass range. Identification of PBH mass in this range may become possible via associated strong primordial gravitational wave detection. There may even be a possibility of detecting \(e^\pm\) lepton asymmetry during \(\sim 1\) year observation of evaporation, if PBH's of a larger mass \(\sim 10^9\) gr now evaporate.

In summary, the inflaton sliding, required for dynamical relaxation towards the zero cosmological constant, gives an ideal framework necessary to generate a finite chemical potential for lepton asymmetry. The simplest scheme uses lepton number violating resonant scattering that occurs in thermal medium. The B–L violation works both in thermal lepto-genesis and primordial black hole evaporation. In either of these scenarios the favored mass scale of heavy Majorana leptons is of order, \(10^{15} \sim 10^{17}\) GeV. Our mechanism does not require the presence of CP violating (CPV) phases in physics beyond the standard model, and can be realized by a minimum extension of standard model incorporating the seesaw mechanism of small neutrino masses. Thus, one neither needs to experimentally detect CPV phases in future neutrino oscillation experiments, nor CPV phases in search for electric dipole moments of various elementary objects are necessary to identify. It is appealing that we already know particle physics necessary to explain matter-antimatter imbalance of our universe, if the unified cosmology is described by multiple scalar-tensor gravity of the sort [8].

If the scenario of primordial black hole evaporation is the origin of matter-antimatter imbalance, it might be possible to determine lepton number violating heavy Majorana masses from footprints of PBH evaporation. With an extra assumption for simplified forms of two \(3 \times 3\) mass matrices, the Majorana \(\mathcal{M}\) and the Dirac mixing \(m\), one might be able to predict the amount of baryon asymmetry using already known neutrino oscillation data. These attempts are left to future works. Nevertheless, it would be welcome if one experimentally finds baryon number violation along with broken B–L number, since the baryo-genesis works equally well as lepto-genesis discussed in the present work. The discovery of proton decay certainly extends our view on the distant future of the universe.
A Thermal average of scattering rate including the region of resonance formation

In the energy region of $s = (p_l + p_H)^2 \geq M_N^2$ the most important is the possibility of resonance production of heavy Majorana lepton of mass $M_N$. We would like to estimate thermally averaged rate of scattering including this energy region. In order to focus on the resonance production, we neglect much less important $u-$channel exchange contribution in (4.2).

The squared amplitude after spin-component summation and average we consider is

$$|A_L|^2 = y^4_{HN} M_N^2 \frac{p_l \cdot \bar{p}_l}{2 E_l E_l} \frac{1}{(s - M_N^2)^2 + \gamma_N^2 M_N^2 / 2}. \quad (A.1)$$

Here $\gamma_N$ is the decay width of $N_R$. In the relevant high energy region of $T \gg m_l, m_H$ it is legitimate to take $s \sim 2 p_l \cdot p_H$; inner product of four-momenta of initial lepton and Higgs boson. Since the thermal distribution functions of bosons and fermions are functions of their energies, we may ignore three-momenta correlation among $\vec{p}_l, \vec{p}_{\bar{l}}, \vec{p}_H$. This leads to an approximate formula of cross section in the free space,

$$|A_L|^2 \approx \frac{1}{2} y^4_{HN} M_N^2 \frac{1}{(2 E_l E_l - M_N^2)^2 + \gamma_N^2 M_N^2 / 2}. \quad (A.2)$$

The thermal average can be done taking the zero chemical potential to a good approximation, hence one should calculate the thermally averaged cross section $\sigma$ as

$$\sigma = \frac{1}{n_l n_H} \int \frac{d^3 p_l d^3 p_H}{(2 \pi)^6} |A_L|^2 f_B \left( \frac{E_H}{T} \right) f_F \left( \frac{E_l}{T} \right), \quad (A.3)$$

where $f_B(x) = 1/(e^x - 1), f_F(y) = 1/(e^y + 1)$ and $n_l = \zeta(3) T^3 / \pi^2, n_H = 3 \zeta(3) T^2 / 4 \pi^2$ are number densities of massless fermions and bosons. This contains, after scaling with the temperature, an integral

$$I \left( \frac{T}{M_N} \right) = \int_0^\infty dx \int_0^\infty dy \frac{1}{(2 x y - M_N^2 / T^2)^2 + \gamma_N^2 M_N^2 / (2 T^4)} \frac{x^2 y^2}{(e^x + 1)(e^y - 1)}, \quad (A.4)$$

times $T^{-4}$. The averaged cross section is given by

$$\sigma \approx \frac{y^4_{HN} M_N^2}{8 \pi^4 T^4} I \left( \frac{T}{M_N} \right). \quad (A.5)$$

The function $I(u)/u^4$ is illustrated in figure 3. One may estimate the resonance temperature $T_R$, taking the average energies for $x$ and $y$, which gives in the narrow width limit

$$T_R \approx \frac{30 \zeta(3)}{\pi^4} \sqrt{\frac{6}{\pi}} M_N \sim 0.34 M_N. \quad (A.6)$$

This gives a good approximation of what we found numerically for the resonance position, as in figure 3. The magnitude of cross section at the resonance pole is $\propto 1 / \gamma_N$, a sensitive quantity to the decay width. The resonance formation may be regarded as a two-step process; the first process being Majorana $N_R$ production via inverse decay and the second its subsequent decay. The enhanced rate $\propto 1 / \gamma_N \propto$ lifetime may be understood by a long lifetime of Majorana particle.
Resonance scattering rate

![Graph showing resonance scattering rate vs temperature/Mn](image)

**Figure 3.** Lepton number violating scattering rate in thermal medium in arbitrary unit. The function $I(x)/x^4$ of (A.4) vs $x = T/M_N$ for three choices of resonance width are illustrated: $\gamma_N/M_N = 0.1$ in solid black, 0.05 in dashed orange, and 0.01 in dash-dotted blue.

The low temperature limit of $I(u)/u^4$ is $\rightarrow 4.3$ as $u \rightarrow 0$, while its high temperature $u \rightarrow \infty$ limit, or $T \gg M_N$, gives

$$\frac{2.404 \ln 2}{4} \frac{1}{u^4} \sim 0.4166 \frac{1}{u^4}. \quad (A.7)$$

The thermally averaged cross sections thus behave in these limits

$$\sigma \rightarrow 4.3 \frac{y_{HN}^4}{8 \pi^4 M_N^3}, \quad \text{as } \frac{T}{M_N} \rightarrow 0, \quad \text{(A.8)}$$

$$\rightarrow 0.42 \frac{y_{HN}^4 M_N^2}{8 \pi^4 T^4}, \quad \text{as } \frac{T}{M_N} \rightarrow \infty. \quad \text{(A.9)}$$

In the resonance region the integral is approximated by

$$I(u) \approx \frac{2\pi \sqrt{2} T^2}{\gamma_N M_N} \int_0^\infty dx \int_0^\infty dy \frac{(xy)^2}{(e^x + 1)(e^y - 1)} \delta \left(2xy - \frac{M_N^2}{T^2}\right) \approx 2.4 \frac{\sqrt{2} \pi^3}{12} \frac{T^2}{\gamma_N M_N}. \quad (A.10)$$

The thermally averaged cross section in the resonance region is given by

$$\sigma \approx 8.77 \frac{y_{HN}^4 M_N^2}{8 \pi^4 T^2} \frac{1}{\gamma_N M_N}, \quad \text{at } T = O \left(\frac{M_N}{3}\right), \quad \text{(A.11)}$$

in the narrow width limit $\gamma_N \ll M_N$.

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