Cosmological Constraints on Gravitino LSP Scenario with Sneutrino NLSP

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Abstract

We study the scenario where a sneutrino is the next lightest supersymmetric particle (NLSP) and decays into a gravitino and standard model particles. The daughter particles such as neutrinos and quarks cause electro- and hadronic showers and affect big-bang nucleosynthesis. It is found that despite a small branching ratio, four-body processes including quarks in the final state give the most stringent constraint on the sneutrino abundance for gravitino mass of $1 - 100$ GeV. Pion production by high energy neutrinos is important when the sneutrinos decay at $\sim 1$ sec. We also discuss the thermal leptogenesis in the sneutrino NLSP scenario.
1 Introduction

Among various possibilities, supersymmetry (SUSY) is a prominent candidate for physics beyond the standard model since it possibly solves various serious problems in particle physics and cosmology. In particular, from cosmology point of view, supersymmetric models contain a good candidate for dark matter; with $R$-parity conservation, the lightest superparticle (LSP) becomes stable and can be dark matter. It is widely known that, if parameters are properly chosen, relic density of the LSP agrees with the dark matter density suggested by WMAP \[1\]

$$\Omega_{CDM} h^2 = 0.105^{+0.007}_{-0.013},$$  \hspace{1cm} (1.1)

where $h$ is the Hubble constant in units of 100 km/sec/Mpc. This fact, as well as other motivations of supersymmetry, provides a strong motivation to consider supersymmetry as a new physics beyond the standard model.

If we consider supersymmetric models, however, several problems may also arise. In particular, in (local) supersymmetric models, superpartner of the graviton, i.e., gravitino, exists. Gravitino is a very weakly interacting particle and it may cause serious cosmological problems \[2\]. In this paper, we concentrate on the case where the gravitino is the LSP.\#1 In this case, gravitino becomes a potential candidate for dark matter.

If gravitino is the LSP, the next to the lightest superparticle (NLSP), which is assumed to be the lightest superparticle in the minimal supersymmetric standard model (MSSM) sector, decays into gravitino with very long lifetime.\#2 Importantly, the lifetime may be longer than 1 sec. Thus, in the evolution of the universe, relic NLSPs may decay during and/or after the big-bang nucleosynthesis (BBN). With the decay of the NLSP, electromagnetic and hadronic showers may be induced. Energetic particles produced in the showers induce photo- and hadro-dissociation and production processes of light elements ($D, ^3He, ^4He, ^6Li$, and $^7Li$). Such processes may significantly change the prediction of standard BBN scenario and, consequently, resultant abundances of light elements may significantly conflict with observations \[10\]. Recent detailed analysis have shown that the model is seriously constrained when the NLSP is neutralino or stau \[11\] \[12\].

However, the lightest neutralino and stau are not the only possible candidates for the lightest superparticle in the MSSM sector. In particular, if the scalar neutrino becomes lightest among the superparticles in the MSSM sector, we expect that the constraints from BBN are drastically relaxed. This is because the dominant decay mode of the parent particle (i.e., sneutrino) is $\tilde{\nu} \rightarrow \psi \nu$, and hence the daughter particles are both weakly interacting.\#3

\#1For recent study for the cases with unstable gravitino, see, for example, \[3\]. (See, also, \[4\] \[5\] \[6\] \[7\] \[8\] \[9\].)
\#2In this paper, we assume $R$-parity conservation.
\#3One may also consider the case where the sneutrino is the LSP and gravitino is the next LSP. Then, the primordial gravitino decays into neutrino-sneutrino pair. In this case, however, gravitino is likely to decay also into charged particles since there should exist charged slepton which is almost degenerate with the sneutrino. Thus, the most stringent constraints from BBN is expected to be from photo-dissociation processes induced by emitted charged particles; constraints on such cases can be read off from old studies. (See, for example, \[6\] \[7\] \[8\].) Thus, in this paper, we concentrate on the case where the gravitino is the LSP.
In order to discuss the BBN reaction in this class of set up, it is crucial to understand the effects of high-energy-neutrino injection on BBN. Such a study was partly done in [13]; in [13], however, only photo-dissociation processes induced by high-energy neutrino injection was discussed, and hadro-dissociation processes were not considered. In addition, the sneutrino NSLP was studied in [11] but the processes induced by the high energy neutrinos were ignored.

In this paper, we consider cosmological constraints on models where (i) the scalar neutrino becomes lightest among the superparticles in the MSSM sector, and (ii) gravitino is the LSP. We pay particular attention to constraints from BBN. Organization of this paper is as follows. In the next section, we discuss properties of sneutrino which plays important role in our scenario. In Section 3 we discuss effects of long-lived sneutrino on BBN. Then, in Section 4 we show results of our numerical study. Implication of our results for thermal leptogenesis scenario is discussed in Section 5. Section 6 is devoted for conclusions and discussion.

2 Properties of Sneutrino

Before discussing cosmology, we first introduce the model we consider. In particular, we discuss properties of sneutrino which plays important role in our study.

Here, we adopt the particle content of the MSSM. Important assumptions in our study are that the lightest superparticle in the MSSM sector is sneutrino \( \tilde{\nu} \), which is in \( SU(2)_L \) doublet slepton \( \tilde{L} \) and that the lightest superparticle is gravitino \( \psi \).

We first discuss when sneutrino can become the lightest superparticle in the MSSM sector. For this purpose, for simplicity, we adopt \( SU(5) \) grand unified model, which is strongly motivated by the fact that, with the particle content of MSSM, three gauge coupling constants meet with a good precision at the grand-unified-theory (GUT) scale \( M_{GUT} \approx 2 \times 10^{16} \) GeV. Then, left-handed (s)leptons and right-handed down-type (s)quarks are in \( \bar{5} \) representations of \( SU(5) \) while other (s)fermions are in \( 10 \) representations.

Detailed mass spectrum of superparticles depends on the mechanism of SUSY breaking. In the case of gauge-mediated SUSY breaking [15], mass spectrum of MSSM superparticles are so constrained as far as the messenger sector respects \( SU(5) \). Consequently, sneutrino does not become the LSP in the MSSM sector.

If we consider supergravity-mediated SUSY breaking scenario, on the contrary, soft SUSY breaking masses for sfermions arise from (non-renormalizable) interactions in Kähler potential. Consequently, in this case, there exist large number of free parameters in the soft SUSY breaking terms. Thus, hereafter, we concentrate on supergravity-mediated SUSY breaking scenario.

\(^{#4} \)In anomaly mediated SUSY breaking scenario [14], gravitino becomes much heavier than superparticles in the MSSM sector and hence it cannot be the LSP. Thus, in anomaly-mediated models, the mass spectrum which we consider cannot be realized and hence our following analysis does not have relevance in such cases. Thus, in the following, we will not consider anomaly-mediated models.
When the soft SUSY breaking terms arise from Kähler interaction, soft SUSY breaking masses for $\tilde{5}$ and $\tilde{10}$ can be independent. In this case, all the sfermions in $\tilde{10}$ representation can become heavier than sneutrino by pushing up the SUSY breaking masses for $\tilde{10}$. In addition, it is usually the case that right-handed down-type squark becomes heavier than left-handed slepton because of the renormalization group effects due to gluino mass, even though their SUSY breaking masses are degenerate at the GUT scale. Thus, it can be easily realized that left-handed slepton becomes the lightest among the sfermions in supergravity-mediated SUSY breaking scenarios. In “conventional” scenarios, however, it is usually the case that the bino mass becomes smaller than the sneutrino mass. This is due to renormalization group effects. Denoting the unified gaugino mass as $M_{1/2}$ and SUSY breaking mass squared of $\tilde{5}$ representation as $\tilde{m}_{\tilde{5}}^2$, electroweak-scale value of the mass squared of $SU(2)_L$ doublet slepton is estimated as 
\[ m_{\tilde{L}}^2(M_{\text{weak}}) = \tilde{m}_{\tilde{5}}^2 + 0.53 M_{1/2}^2. \]
(Here, $M_{1/2}$ and $\tilde{m}_{\tilde{5}}^2$ are defined at the GUT scale.) In addition, bino mass is given by 
\[ M_1(M_{\text{weak}}) \simeq 0.41 M_{1/2}. \]
Thus, with the assumption of $\tilde{m}_{\tilde{5}}^2 \geq 0$, $M_1^2$ becomes smaller than $\tilde{m}_{\tilde{L}}^2(M_{\text{weak}})$ and sneutrino becomes heavier than bino.

However, we can easily avoid this conclusion by assuming $\tilde{m}_{\tilde{5}}^2 < 0$. As we mentioned, some part of SUSY breaking parameters originate from direct interaction between SUSY breaking fields and observable-sector fields. Denoting the superfields for SUSY breaking field and observable-sector field as $\hat{z}$ and $\hat{\phi}$, respectively, the following interaction may exist:
\[ L_{\text{int}} = \frac{\lambda}{M_{\text{Pl}}} \int d^4 \theta \hat{z}^* \hat{z} \hat{\phi}^* \hat{\phi}, \]
where $\lambda$ is a coupling constant. After SUSY is broken, scalar component in $\hat{\phi}$ acquires SUSY breaking mass squared as $-\frac{\lambda}{M_{\text{Pl}}} |F_z|^2$, where $F_z$ is the $F$-component of $\hat{z}$. Assuming $\lambda > 0$, negative contribution to mass squared of scalar field is obtained. Importantly, $\lambda$ is a free parameter in Kähler potential and its sign is unknown from low-energy effective theory point of view. Thus, we assume that the $\lambda$ parameter for doublet slepton is positive and that $\tilde{m}_{\tilde{L}}^2(M_{\text{weak}}) < M_1^2$. Of course, in this case, there exists true minimum where the doublet slepton acquires non-vanishing vacuum expectation value. This is not a serious problem since the transition rate to the true minimum is small enough so that the transition to the true minimum does not occur in the cosmic time scale.

When the relation $\tilde{m}_{\tilde{L}}^2(M_{\text{weak}}) < M_1^2$ holds, sneutrino may become the lightest superparticle in the MSSM sector in large fraction of the parameter space. Indeed, neglecting left-right mixing, masses of sneutrino $\tilde{\nu}$ and left-handed charged lepton $\tilde{l}_L$ are given by
\[ m_{\tilde{\nu}}^2 = \tilde{m}_{\tilde{L}}^2(M_{\text{weak}}) + \frac{1}{2} m_Z^2 \cos 2\beta, \]
(Here, we have neglected the effect of Yukawa coupling constants; it may slightly decrease $\tilde{m}_{\tilde{L}}^2(M_{\text{weak}})$.)
\[ m_{i_L}^2 = \tilde{m}_{\tilde{l}}^2(M_{\text{weak}}) + \left( \sin^2 \theta_W - \frac{1}{2} \right) m_Z^2 \cos 2\beta, \]  

(2.5)

where \( \theta_W \) is the Weinberg angle, and \( \tan \beta \) is the ratio of vacuum expectation values of two Higgs bosons. Thus, when \( \tan \beta > 1 \), \( \tilde{\nu} \) becomes lighter than \( \tilde{l}_L \). Thus, hereafter, we consider the case where the sneutrino is the lightest superparticle in the MSSM sector. In addition, as we mentioned, we consider the case where the LSP is the gravitino \( \psi \). In this class of set up, BBN may be affected by late-time decay of primordial sneutrino which freezes out from the thermal bath when the cosmic temperature drops below the sneutrino mass.

Next, let us consider decay processes of \( \tilde{\nu} \). Decay of the sneutrino is dominated by the two-body process \( \tilde{\nu} \rightarrow \psi \nu \); as we will see, decay rates for three- and four-body decay processes are much smaller than \( \Gamma_{\tilde{\nu} \rightarrow \psi \nu} \). Thus, the decay rate of sneutrino is given by

\[
\Gamma_{\tilde{\nu}} \simeq \Gamma_{\tilde{\nu} \rightarrow \psi \nu} = \frac{m_{\tilde{\nu}}^5}{48\pi m_{3/2}^2 M_s^2} \left( 1 - \frac{m_{3/2}^2}{m_{\tilde{\nu}}^2} \right)^4,
\]  

(2.6)

where \( m_{3/2} \) and \( m_{\tilde{\nu}} \) are the masses of gravitino and sneutrino, respectively, while \( M_s \approx 2.4 \times 10^{18} \) GeV is the reduced Planck scale. The lifetime of sneutrino \( \tau_{\tilde{\nu}} \) is calculated by using the two body decay rate given in Eq. (2.6). In Fig. 1 we plot contours of constant \( \tau_{\tilde{\nu}} \) on \( m_{3/2} \) vs. \( m_{\tilde{\nu}} \) plane. Importantly, as the gravitino mass becomes smaller, the decay rate of \( \tilde{\nu} \) becomes larger since the interaction with the longitudinal component of the gravitino is more enhanced. Thus, if the gravitino mass becomes smaller than \( \sim 0.1 - 1 \) GeV, lifetime of sneutrino becomes shorter than \( \sim 1 \) sec. In this case, primordial sneutinos in early universe decay before BBN starts and no constraints are obtained from BBN. On the contrary, with larger gravitino mass, lifetime becomes longer and we should carefully consider hadro- and photo-dissociation processes induced by the decay of sneutrino.

Even though the dominant decay of \( \tilde{\nu} \) is via two-body process, it is also necessary to consider three- and four-body decay processes in the study of non-standard BBN reactions. In particular, processes like \( \tilde{\nu} \rightarrow \psi \nu Z^{(*)} \) and \( \tilde{\nu} \rightarrow \psi l W^{(*)} \), followed by \( Z^{(*)} \rightarrow f \bar{f} \) and \( W^{(*)} \rightarrow f \bar{f} \), are important. (Here, the superscript “*” is for vertual particles. In addition, \( f \) represents quarks and leptons, and \( l \) denotes charged lepton.) This is because charged and colored particles are produced only via three- and four-body decay processes. Those charged and colored particles efficiently interact with background particles and induce photo- and hadro-dissociation processes. As we will see in the following sections, these three- and four-body processes may significantly change predictions of standard BBN scenario in some region of parameter space.

In order to study effects of decay modes with colored and/or charge particles in the final state, we calculate the decays rate for the processes \( \tilde{\nu} \rightarrow \psi \nu f \bar{f} \), \( \tilde{\nu} \rightarrow \psi l f \bar{f} \), and \( \tilde{\nu} \rightarrow \psi \nu \gamma \). Here, we approximated that the neutralinos and charginos are purely gaugino (or Higgsino). In addition, we neglect effects of Yukawa coupling constants. Then, for the processes \( \tilde{\nu} \rightarrow \psi \nu f \bar{f} \) and \( \tilde{\nu} \rightarrow \psi l f \bar{f} \), we take account of the Feynman diagrams shown in Fig. 2. (In our analysis, we do not use the narrow-width approximation for the productions of “on-shell” \( Z \) and \( W \)-bosons; effects of “on-shell” \( Z \) and \( W \) productions are taken into account
at the phase-space region where the invariant mass of $f \bar{f}$ (or $f \bar{f}'$) system becomes close to the gauge-boson mass.) In Fig. 3 we plot the contours of constant hadronic branching ratio, which is defined as

$$B_h = \frac{1}{\Gamma_{\tilde{\nu}}} \left[ \sum_q \Gamma_{\tilde{\nu} \to \psi \nu qq} + \sum_q \Gamma_{\tilde{\nu} \to \psi \nu q' q''} \right].$$

(2.7)

As one can see, the hadronic branching ratio is very small. As we will see, however, hadronic decay processes play important role in discussing dissociation of light elements. We also calculate decay rate for the process $\tilde{\nu} \to \psi \nu \gamma$, which occurs with the diagram like Fig. 2(c) (without attaching $f \bar{f}$ final state). We found, however, that the decay rate for this process is negligibly small with the parameter set we use in our numerical analysis.

In order to study effects of photo-dissociation processes, we also calculate averaged “visible energy” emitted from the decay of one sneutrino:

$$E_{\text{vis}} = \frac{1}{\Gamma_{\tilde{\nu}}} \left[ \langle E_{\text{vis}}^{\tilde{\nu} \to \psi \nu \gamma} \rangle \Gamma_{\tilde{\nu} \to \psi \nu \gamma} + \sum_f \langle E_{\text{vis}}^{\tilde{\nu} \to \psi \nu f f'} \rangle \Gamma_{\tilde{\nu} \to \psi \nu f f'} + \sum_{f'} \langle E_{\text{vis}}^{\tilde{\nu} \to \psi \nu f f'} \rangle \Gamma_{\tilde{\nu} \to \psi \nu f f'} \right],$$

(2.8)

where $\langle E_{\text{vis}}^{\tilde{\nu} \to \cdots} \rangle$ are averaged energy carried away by charged particles and photon in individual decay modes $\tilde{\nu} \to \cdots$. We have checked that the quantity $\langle E_{\text{vis}}^{\tilde{\nu} \to \cdots} \rangle$ are typically $O(10\%)$ of the sneutrino mass when $m_{3/2} \ll m_{\tilde{\nu}}$. However, $E_{\text{vis}}$ is much smaller than $m_{\tilde{\nu}}$ since the branching ratios for events with charged particle(s) in the final state are small.

Figure 1: Contours of constant lifetime of the sneutrino on $m_{3/2}$ vs. $m_{\tilde{\nu}}$ plane. We have shaded the region with $m_{\tilde{\nu}} < m_{3/2}$, which we are not interested in.
3 BBN with Long-Lived Sneutrino

Now, we are at the position to discuss effects of long-lived sneutrino on BBN. As we mentioned, there are two types of decay modes of sneutrino; one is two-body decay mode while the other are three- and/or four-body ones. Reactions caused by these decay modes are different, so we discuss effects of these decay modes separately.

3.1 Two-body decay

First, we discuss effects of the dominant decay process $\tilde{\nu} \rightarrow \psi \nu f \bar{f}$. The gravitino produced in the decay is a very weakly interacting particle, so it is irrelevant for BBN.

Neutrino, on the contrary, may affect abundances of light elements. Once emitted, the energetic neutrinos may scatter off background leptons via weak interaction. Consequently, several kinds of particles may be pair-produced.

First, charged leptons may be produced via the following processes:

$$\nu_i \bar{\nu}_{i,\text{BG}} \rightarrow e^- e^+,$$

Figure 2: Feynman diagrams for the decay processes $\tilde{\nu} \rightarrow \psi \nu f \bar{f}$ ((a) − (d)) and $\tilde{\nu} \rightarrow \psi l f \bar{f}'$ ((e) − (h)). Here, $\tilde{B}$, $\tilde{W}^0$, and $\tilde{W}^\pm$ represent bino, neutral wino, and charged wino, respectively.
Figure 3: Contours of constant hadronic branching ratio $B_h$ on $m_{3/2}$ vs. $m_{\nu}$ plane.

$\nu_i\bar{\nu}_i,\text{BG} \to \mu^-\mu^+,$ (3.2)
$\nu_{\mu}\bar{\nu}_{\mu},\text{BG} \to \mu^-e^+,$ (3.3)
$\nu_{\tau}\bar{\nu}_{\tau},\text{BG} \to \mu^-\mu^+,$ (3.4)

where $i=e,\mu,\tau$ is flavor index, and the subscript “BG” is for background particles. The muons emitted in the above processes quickly decay into electrons and neutrinos. Thus, the above processes produce energetic electrons (and positrons) which cause electromagnetic cascade and energetic photons in the cascade induce photo-dissociation processes of light elements. Effects of these processes have been already studied in [13]; it was pointed out that, in some part of the parameter space, abundances of light elements are significantly affected if energetic neutrinos are injected.

Other possible effect is due to the production of pion pair which was not considered in [13]. High energy neutrinos scatter off the background neutrinos and electrons and produce pions as

$\nu_i\bar{\nu}_i,\text{BG} \to \pi^-\pi^+,$ (3.5)
$\nu_{\tau}\bar{\nu}_{\tau},\text{BG} \to \pi^0\pi^-.$ (3.6)

The nucleus-pion interaction rate is $\sim 10^8 \text{ sec}^{-1} \times (T/\text{MeV})^3$ which is larger than the decay rate of the charged pion ($\sim 4 \times 10^7 \text{ sec}^{-1}$) for $T \lesssim 1 \text{ MeV}$. Therefore, the charged pions produced at $T \sim 1 \text{ MeV}$ scatter off the background nuclei and change protons (neutrons) into neutrons (protons) via

$\pi^-p \to n\pi^0 \text{ or } n\gamma,$ (3.7)
$\pi^+n \to p\pi^0 \text{ or } p\gamma,$ (3.8)
which increases the n-p ratio and synthesizes more $^4\text{He}$. On the other hand, because of very short lifetime, the neutral pions decay before they scatter off the background nuclei.

In order to estimate effects of the high energy neutrino induced processes, we have to numerically solve the Boltzmann equation describing the time evolution of the high energy neutrino spectrum taking account of neutrino-neutrino (charge lepton) scattering as well as production of charged leptons and pions. (The details of the Boltzmann equation and its solution will be presented elsewhere [19].)

### 3.2 Three- and four-body processes

Even though the branching ratio for the three- and four-body decay processes are much smaller than 1, such decay processes are very important since colored and/or charged particles are directly emitted from these decay processes. Energetic colored and charged particles may significantly change the prediction of standard BBN scenario. Effects of these colored and charged particles are classified into three categories: photo-dissociations, hadro-dissociations, and $p \leftrightarrow n$ conversion.

Photo-dissociation processes are induced by energetic photons in electromagnetic shower which is caused by charged particles emitted from $\bar{\nu}$. With given background temperature, the distribution function of energetic photons depends on total amount of energy injected by particles with electromagnetic interaction, and is insensitive to the energy spectrum of primary particles. Thus, once $E_{\nu\bar{\nu}}$ is obtained, the energy distribution of energetic photons in the electromagnetic shower can be obtained. Then, photo-dissociation rates are obtained by convoluting energy distribution function and cross sections of photo-dissociation reactions. For details of our treatment of photo-dissociation processes, see [20].

For the study of hadro-dissociation processes, it is necessary to obtain energy distributions of (primary) hadrons which are produced after the hadronization of quarks emitted from $\bar{\nu}$. As a first step to derive distribution functions of hadrons, we calculate invariant-mass distribution of $q\bar{q}^{(i)}$ system produced by $\bar{\nu} \rightarrow \psi\nu q\bar{q}$ and $\bar{\nu} \rightarrow \psi lqq'$. For example, for the neutral-current event $\bar{\nu} \rightarrow \psi\nu q\bar{q}$, we numerically estimate the following quantity for each quark flavor $q$:

$$
\frac{d\Gamma_{\bar{\nu} \rightarrow \psi\nu q\bar{q}}}{dm_{q\bar{q}}^2} = \frac{N_c}{8\pi^2m_{\bar{\nu}}} \int dm_{\nu\nu q\bar{q}}^2 d\Phi_{\psi,(\nu q\bar{q})} d\Phi_{\psi,(qq')}\left|\mathcal{M}_{\bar{\nu} \rightarrow \psi\nu q\bar{q}}\right|^2, \quad (3.9)
$$

where $N_c = 3$ is the color factor, $\mathcal{M}_{\bar{\nu} \rightarrow \psi\nu q\bar{q}}$ is the matrix element, and $m_{\nu\nu q\bar{q}}$ and $m_{\nu q\bar{q}}$ are invariant masses of $q\bar{q}$ and $\nu q\bar{q}$ system, respectively. In addition, $d\Phi_{x,y}$ represents (infinitesimal) two-body phase space of $x$ and $y$. For each quark flavor $q$, spectra of hadrons are calculated by PYTHIA code [21]. We denote the spectrum of nucleus $N$ in the center-of-mass frame of $q\bar{q}$, which is obtained after the hadronization of one pair of $q\bar{q}$, as

$$
f_N^{(\text{cm})}(E_N^{(\text{cm})}, m_{q\bar{q}}) = \left[\frac{dN_N}{dE_N^{(\text{cm})}}\right]_{m_{q\bar{q}}}, \quad (3.10)
$$
where $N_N$ denotes the number of nucleus $N$. It should be noted that $f_N^{(cm)}$ is not the spectrum in the rest frame of $\tilde{\nu}$. In order to take account of the effect of boost, we numerically calculate the averaged boost factor of the $q\bar{q}$ system for as a function of $m_{q\bar{q}}$:

$$
\tilde{\gamma}_{q\bar{q}}(m_{q\bar{q}}) = \frac{1}{m_{q\bar{q}}} \int \frac{dm_{q\bar{q}}^2}{m_{q\bar{q}}} d\Phi_{\psi, (q\bar{q})} d\Phi_{\nu, (q\bar{q})} d\Phi_{q, \bar{q}} \frac{|M_{\tilde{\nu} \rightarrow \psi q\bar{q}}|^2}{|M_{\tilde{\nu} \rightarrow \psi q\bar{q}}|^2},
$$

(3.11)

where $E_{q\bar{q}}$ is the energy of $q\bar{q}$ system in the rest frame of $\tilde{\nu}$. With this quantity, we estimate the spectrum of the nucleus $N$ in the rest frame of $\tilde{\nu}$, approximating that the boost factor for a fixed value of $m_{q\bar{q}}$ is universally $\tilde{\gamma}_{q\bar{q}}(m_{q\bar{q}})$. In addition, for simplicity, we approximate that the distribution of initial $q$ (and hence $\bar{q}$) jet is isotropic in the center-of-mass frame of $q\bar{q}$. Then, we obtain

$$
\left[\frac{dN_N}{dE_N}\right]_{\tilde{\nu} \rightarrow \psi q\bar{q}} = \frac{1}{2\Gamma_{\tilde{\nu} \rightarrow \psi q\bar{q}}} \int d\cos \theta_q dm_{q\bar{q}}^2 \frac{\partial f_N^{(cm)}}{\partial E_N} \frac{\partial f_N^{(cm)}}{\partial m_{q\bar{q}}} \frac{d\Gamma_{\tilde{\nu} \rightarrow \psi q\bar{q}}}{dm_{q\bar{q}}},
$$

(3.12)

where, in the above formula, $E_N^{(cm)}$ is related to $E_N$ as

$$
E_N = \tilde{\gamma}_{q\bar{q}} \left( E_N^{(cm)} + \tilde{\beta}_{q\bar{q}} \sqrt{E_N^{(cm)2} - m_N^2 \cos \theta_q} \right),
$$

(3.13)

with $\tilde{\beta}_{q\bar{q}} = \sqrt{1 - \tilde{\gamma}_{q\bar{q}}^2}$. Effects of charged-current events $\tilde{\nu} \rightarrow \psi lq\bar{q}'$ are also treated in the same way.

We have calculated the spectra of $p$ and $n$ taking account of all possible neutral- and charged-current events. These hadrons cause hadronic shower and induce hadro-dissociation processes. In our analysis, in addition, we have also estimated the number of charged pions produced by the decay of $\tilde{\nu}$ with the same procedure as the case of $p$ and $n$. Such charged pions become the source of $p \leftrightarrow n$ conversion process, which changes the number of $^4\text{He}$. Once the spectra of hadrons are obtained, effects of hadro-dissociation and $p \leftrightarrow n$ conversion are studied following [7, 8].

## 4 Numerical Results

Now, we are at the position to show our numerical results. In our analysis, we have followed the evolutions of the number densities of light elements by numerically solving Boltzmann equations. For this purpose, we have modified Kawano code [22], including non-standard processes discussed in the previous section:

- $p \leftrightarrow n$ conversion by pions produced via $\nu\nu_{BG} \rightarrow \pi^+\pi^-$ and three- and four-body decays of $\tilde{\nu}$.
- Photo-dissociation processes by charged leptons produced via $\nu\nu_{BG} \rightarrow l^+l^-$. 


• Hadro-dissociation processes by hadrons produced via three- and four-body decay processes of $\tilde{\nu}$.

• Photo-dissociation processes by charged particles produced via three- and four-body decays of $\tilde{\nu}$.

In order to derive typical constraint, in our analysis, the primary neutrino emitted from NLSP sneutrino is treated as equal-weight admixture of $\nu_e$, $\nu_\mu$, and $\nu_\tau$. Notice that the most important constraints, which are from hadro-dissociation processes, do not change even if we adopt different assumption (although the constraints from the two-body decay of $\tilde{\nu}$ may have slight dependence on the flavor of primary neutrino). Then, we compare the results of our calculation with light-element abundances inferred from observations; we calculate $\chi^2$ variable defined in [8] and obtained 95 % C.L. constraints. The observational constraints adopted in this paper are summarized in Appendix A.

Effects of long-lived sneutrino on BBN depends on primordial abundance of sneutrino. We parameterize the primordial abundance is parameterized by yield variable, which is defined as the ratio of number density and total entropy density (at $t \ll \tau_\tilde{\nu}$):

$$Y_\tilde{\nu} \equiv \left[ \frac{n_\tilde{\nu}}{s} \right]_{t \ll \tau_\tilde{\nu}} .$$

With this quantity, the number density of sneutrino is given by $n_\tilde{\nu} = sY_\tilde{\nu}e^{-t/\tau_\tilde{\nu}}$. Importantly, $Y_\tilde{\nu}$ depends on thermal history of the universe, so we consider several cases.

We first consider the case where the sneutrinos are thermal relics; in this case, we assume "standard" evolution of the universe when the cosmic temperature is below the mass scale of SUSY particles $m_{\text{SUSY}}$ which is taken to be 100 GeV – 1 TeV. When the cosmic temperature $T$ is higher than $m_{\text{SUSY}}$, all the MSSM particles are thermalized and their number density is of order $T^3$. As the temperature becomes lower than $m_{\text{SUSY}}$, on the contrary, density of MSSM superparticles are Boltzmann-suppressed and the number densities of most of the superparticles become negligibly small as $T \to 0$. Only the exception is the number density of the lightest superparticle in the MSSM sector, in our case, sneutrino. With $R$-parity conservation, $\tilde{\nu}$ decays only into gravitino and some other particle(s). Chemical equilibrium of sneutrino is maintained by pair annihilation process, whose rate becomes smaller than the expansion rate of the universe at some epoch. After this epoch, sneutrino freezes out from the thermal bath and its number density in the comoving volume is (almost) conserved at the cosmic time when $t \ll \tau_\tilde{\nu}$. Relic density of sneutrino in such a case was studied in [17], and the yield variable is given by

$$Y_\tilde{\nu} \simeq 2 \times 10^{-14} \times \left( \frac{m_\tilde{\nu}}{100 \text{ GeV}} \right).$$

Using the yield variable given above, we calculate the abundances of light elements for fixed values of $m_{3/2}$ and $m_\tilde{\nu}$. Then, comparing the results of our calculation with observational constraints on the primordial abundances of light elements, we derive constraints on $m_{3/2}$ vs. $m_\tilde{\nu}$ plane.
Figure 4: Constraints on $m_{3/2}$ vs. $m_{\tilde{\nu}}$ plane when the relic sneutrinos have thermal origin. Regions inside the contours are excluded.

In Fig. 4, we show constraints on $m_{3/2}$ vs. $m_{\tilde{\nu}}$ plane when the relic sneutrinos have thermal origin. In this case, we have found that the most important constraints are obtained from the overproductions of D and $^6\text{Li}$. In order to understand the behavior of the constraints, it should be noticed that the lifetime $\tau_{\tilde{\nu}}$ becomes shorter as the gravitino mass becomes smaller. Consequently, when the gravitino mass is small enough, sneutrino decays at very early stage of BBN and no constraint is obtained. On the contrary, when $10^2 \text{ sec} \lesssim \tau_{\tilde{\nu}} \lesssim 10^7 \text{ sec}$, the background $^4\text{He}$ (which we denote $\alpha_{BG}$) is effectively dissociated by the energetic hadrons produced in the hadronic shower. In this case, overproduction of D may occur as a result of hadro-dissociation of $\alpha_{BG}$. In addition, energetic T and $^3\text{He}$ is also produced by the hadro-dissociation process. Such T and $^3\text{He}$ become sources of non-thermal production of $^6\text{Li}$ via $T + \alpha_{BG} \to ^6\text{Li} + n$ and $^3\text{He} + \alpha_{BG} \to ^6\text{Li} + p$. As the lifetime becomes longer, the energetic hadrons are stopped by the scattering processes with background particles and hence the effects of hadro-dissociations become inefficient. In addition, we also have found that effects of photo-dissociation are not important in this case. The constraint obtained here is qualitatively consistent with the result in [11], but our constraint extends to the region with smaller gravitino mass.

So far, we have considered the case where the relic sneutrinos have thermal origin; in such a case, the relic density of $\tilde{\nu}$ is given by Eq. (4.2). In non-standard cases, however, relic density of $\tilde{\nu}$ may become larger. For example, for the case where late-decaying scalar condensation exists, superparticles may be produced at the decay time of scalar condensation. If the reheating temperature after the decay of scalar condensation is lower than $\sim 10 \text{ GeV}$, relic density of sneutrino becomes larger than that given in Eq. (4.2). Thus, taking $Y_{\tilde{\nu}}$...
as a free parameter, we calculate the abundances of light elements.

In Fig. 5 we show constraints on $m_{3/2}$ vs. $m_\tilde{\nu} Y_{\tilde{\nu}}$ plane. Here, we take $m_\tilde{\nu} = 300$ GeV; with this value of sneutrino mass, the light-element abundances are consistent with observational constraints if the relic sneutrinos have thermal origin. However, as we adopt larger value of $Y_{\tilde{\nu}}$, some region of the parameter space is excluded. As one can see in Fig. 5 the most important constraints are from overproduction of $^4\text{He}$ when $m_{3/2} \lesssim 0.4$ GeV and from overproductions of $^3\text{He}$ and $^6\text{Li}$ when $m_{3/2} \gtrsim 0.4$ GeV. In addition, when the gravitino mass becomes close to the sneutrino mass, the lifetime of the sneutrino becomes relatively long and photo-dissociation of $^{12}\text{C}$ becomes effective. In this case, overproduction of $^3\text{He}$ may also occur. On the other hand, for small gravitino mass ($\lesssim 0.4$ GeV) the lifetime of the sneutrino is short ($\lesssim 40$ sec) and pion production by the high energy neutrino emitted in the two-body decay becomes significant, which leads to change $p-n$ ratio and $^4\text{He}$ overproduction.

Gravitinos are produced by the decay of sneutrinos. Since gravitino is stable, they contribute to the present mass density of the universe. Density parameter of the gravitino produced by the sneutrino decay is obtained by

$$\Omega_{3/2} = \frac{m_{3/2} s_{\text{bow}} Y_{\tilde{\nu}}}{\rho_c},$$

(4.3)
where \( s_{\text{now}} \) is the entropy density at the present universe and \( \rho_c \) is the critical density. Requiring \( \Omega_{\beta/2} < \Omega_{\text{CDM}} \), \( Y_{\beta} \) is constrained from above. The bound is also shown in Fig. 5.\(^6\)

It should be noted that, when \( m_{3/2} \gtrsim 0.1 \) GeV, BBN constraints on \( Y_{\beta} \) are more stringent than that obtained from the overproduction of the gravitino in most of the cases.

## 5 Implication for Leptogenesis

In this section, we discuss implication of sneutrino-NSLP scenario for thermal leptogenesis. It is widely known that the present baryon asymmetry of the universe may originate from non-equilibrium decay of right-handed (s)neutrino which has frozen out from the thermal \[^{23}\].

In order to realize thermal leptogenesis, however, high enough reheating temperature is needed \[^{24}\]. Here, we define reheating temperature as

\[
T_R \equiv \left( \frac{10}{g_* \pi^2} M_*^2 \Gamma_{\text{inf}}^2 \right)^{1/4},
\]

where \( g_* \) is the effective number of massless degrees of freedom. (We take the MSSM value \( g_* = 228.75 \) in our calculation.) Then, detailed calculation based on MSSM shows that \( T_R \gtrsim 1 \times 10^9 \) GeV is required in order to generate large enough baryon asymmetry \[^{25}\].

If the reheating temperature is as high as \( \sim 10^9 \) GeV, however, large number of gravitinos are produced by scattering processes of MSSM particles in thermal bath. Such gravitinos usually cause serious cosmological problems. If gravitino is unstable, various colored and/or charged particles are produced when gravitinos decay, resulting in serious dissociations of light elements during/after BBN. If \( T_R \sim 10^9 \) GeV, light-element abundances are too much affected to be consistent with observations unless gravitino mass is extremely large \( (m_{3/2} \gtrsim 10 \) TeV) \[^{8,3}\].

With stable gravitino, on the contrary, the present mass density of gravitino may become too large \[^{10}\]. Since the number density of gravitino is approximately proportional to the reheating temperature, we obtain upper bound on the reheating temperature. This fact severely constrains the scenario of thermal leptogenesis \[^{17}\]. Importantly, as the gravitino mass becomes smaller, production of longitudinal mode of gravitino is enhanced and upper bound on \( T_R \) becomes more stringent.

We have numerically solved Boltzmann equations from the inflaton-dominated epoch to radiation dominated epoch and derived upper bound on the reheating temperature. In our calculation, we have assumed that, in the inflaton-dominated epoch, inflaton potential is well approximated by parabolic potential and that inflaton is rapidly oscillating. Then,

\[^6\]In fact, gravitinos are also produced by scattering processes of particles in thermal bath. For details, see the next section.

\[^7\]Reheating temperature used in \[^{25}\] (which we denote \( T_R^{\text{GNRRS}} \)) is related to our definition as \( T_R^{(\text{GNRRS})} = \sqrt{3} T_R \).
Figure 6: Contours of constant maximal possible reheating temperature after inflation (or any other entropy production) in order to avoid overproduction of gravitino. The horizontal axis is the gravitino mass while the vertical axis is bino mass.

Denoting energy densities of radiation and inflaton as $\rho_{\text{rad}}$ and $\rho_{\text{inf}}$, respectively, and the number density of gravitino as $n_{3/2}$, they obey the following equations

\begin{align}
\frac{d\rho_{\text{rad}}}{dt} &= -4H\rho_{\text{rad}} + \Gamma_{\text{inf}}\rho_{\text{inf}}, \quad (5.2) \\
\frac{d\rho_{\text{inf}}}{dt} &= -3H\rho_{\text{inf}} - \Gamma_{\text{inf}}\rho_{\text{inf}}, \quad (5.3) \\
\frac{dn_{3/2}}{dt} &= -3Hn_{3/2} + \langle \sigma_{\text{tot}} v_{\text{rel}} \rangle \rho_{\text{rad}}, \quad (5.4)
\end{align}

where $H$ is the expansion rate of the universe, and $\rho_{\text{rad}} = \frac{\zeta(3)}{\pi^4}T^3$. In addition, $\langle \sigma_{\text{tot}} v_{\text{rel}} \rangle$ is “thermally averaged” total cross section (times relative velocity), which is given in [26]. When the gravitino mass is small, $\langle \sigma_{\text{tot}} v_{\text{rel}} \rangle$ is proportional to squared of gaugino masses and, consequently, the resultant upper bound is sensitive to the gaugino masses. Here, we adopt GUT relation among gaugino masses

\begin{align}
\frac{M_3}{g_3^2} = \frac{M_2}{g_2^2} = \frac{3}{5} \frac{M_1}{g_1^2}, \quad (5.5)
\end{align}

where $M_3$, $M_2$, $M_1$ and $g_3$, $g_2$, $g_1$ are gaugino masses and gauge coupling constants for $SU(3)_C$, $SU(2)_L$, and $U(1)_Y$ gauge interactions, respectively.

Numerically solving Eqs. (5.2) – (5.4) from the epoch with $H \gg \Gamma_{\text{inf}}$ to the present epoch, we calculate the density parameter of gravitino $\Omega_{3/2} = \frac{m_{3/2} n_{3/2}}{\rho_c}$ as a function of
reheating temperature. In deriving upper bound on $T_R$, we require that $\Omega_{3/2}$ be smaller than the best-fit value of the dark-matter density reported by WMAP (see Eq. 1.1); in other words, the upper bound $T_R^{(\text{max})}$ satisfies the relation $\Omega_{3/2}(T_R^{(\text{max})}) h^2 = 0.105$. Thus, for the case where $T_R = T_R^{(\text{max})}$, gravitino dark matter is realized.

The result is shown in Fig. 6. As we mentioned, the gravitino production cross section is more enhanced with larger value of gaugino mass and smaller value of $m_{3/2}$. Thus, in those limits, upper bound on $T_R$ becomes more stringent. Comparing Fig. 4 with Fig. 6 we can see that reheating temperature of $O(10^9 \text{ GeV})$ is possible without conflicting BBN constraints for $20 \text{ GeV} \lesssim m_{3/2} \lesssim 400 \text{ GeV}$ and $M_1 \lesssim 400 \text{ GeV}$. (Notice that $m_{\tilde{\nu}}$ is smaller than $M_1$ since consider the case where sneutrino is the lightest superparticles in the MSSM sector.) Thus in the scenario with gravitino NLSP and gravitino LSP, thermal leptogenesis scenario is possible in the above-mentioned parameter region. In this region, in addition, it should be also noticed that gravitino is a viable candidate for dark matter.

It is important to note that the upper bound given in Fig. 6 is applicable irrespective of what the NLSP is. Thus, for thermal leptogenesis, gravitino mass should be $O(10 \text{ GeV})$ or larger. This fact gives serious constraint. When bino or charged slepton is the NLSP (with gravitino being the LSP), photo-dissociation processes provide stringent constraints on the parameter region $20 \text{ GeV} \lesssim m_{3/2} \lesssim 400 \text{ GeV}$ and $M_1 \lesssim 400 \text{ GeV}$ where $T_R \gtrsim 10^9 \text{ GeV}$ is allowed [11, 12]. In particular, photo-dissociation processes of $^4\text{He}$ enhance the ratio $^3\text{He}/\text{D}$ too much. Thus, if we adopt upper bound on the ratio $^3\text{He}/\text{D}$ given in Eq. (A.2), such parameter region is excluded. Thus we have to conclude that thermal leptogenesis is difficult for the cases with bino NLSP and charged-slepton NLSP. On the contrary, as we have seen, thermal leptogenesis is possible in large parameter space if we consider sneutrino NLSP. Thus, the scenario with sneutrino NLSP and gravitino LSP has an advantage in realizing thermal leptogenesis.

Before closing this section, we comment on what happens if we relax the GUT relation on the gaugino masses. With the GUT relation, gluino becomes much heavier than bino. Since the gravitino production cross sections are proportional to squared of gaugino masses, primordial gravitinos are mostly produced by processes with gluino in initial and/or final states. If the gluino mass is smaller, gravitino production rate is suppressed for a fixed value of $M_1$. Indeed, in some classes unified models, GUT relation may not hold; models with hypercolor [27] is one of such cases [28]. Then, with smaller gluino mass, upper bound on the reheating temperature can become higher for a fixed value of bino mass. In such case, thermal leptogenesis may be possible even with bino NLSP and stau NLSP.

6 Conclusions

In this paper we have investigated the BBN constraint on the scenario in which the sneutrino is NLSP and decays into a gravitino (i.e., LSP) and the standard model particles. Sneutrino

[In fact, gravitinos are also produce by the decay of sneutrino. However, we neglect such contribution since it is subdominant for most of the cases as as we can estimate from Eq. (1.12).]
mainly decays into a gravitino and a neutrino which scatters off the background neutrinos and electrons producing pions and high energy electrons. Moreover quarks and photons are also produced via three- or four-body decay of the sneutrino although their branching ratios are small. For the case where the sneutrinos are thermal relics, the most stringent constraint comes from overproduction of D and $^6\text{Li}$ produced in the hadron showers which are induced by the four-body decays. We also derived the constraint on the sneutrino abundance assuming the sneutrinos are produced non-thermally. We have found that the BBN constraint is more stringent than that from the cosmic density of the gravitino in the wide range of gravitino mass. It is also found that pion production by high energy neutrinos from the two-body decay becomes important when the sneutrinos decay at $\lesssim 10$ sec.

Since the main decay mode ($\bar{\nu} \rightarrow \psi \nu$) leads to a less stringent constraint, the sneutrino NLSP scenario allows a larger parameter space than other scenarios like stau NLSP. In fact, it has been shown that the thermal leptogenesis is realized. 

Note Added: After finalizing the manuscript, we noticed the paper [29] which has some overlap with our analyses.

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A Observational Constraints on Light Elements

In this appendix, we summarize the observational constraints on light-element abundances which we use in our analysis. For some light elements, we take account of recent developments in observations, so the observational constraints adopted in this paper is different from those used in the old study [7, 8].

The primordial value of the ratio \text{D}/\text{H} is measured in the high red-shift QSO absorption systems. Recently a new deuterium data was obtained from observation of the absorption system at the red-shift $z = 2.70$ towards QSO SDSS1558-0031 [30]. The reported value of the deuterium abundance is log(D/H) = $-4.48 \pm 0.06$. Combined with the previous data [31, 32, 33, 34, 35], it is reported that the primordial abundance is given by

$$ (\text{D}/\text{H})_p = (2.82 \pm 0.26) \times 10^{-5}. $$

(A.1)

For the constraint on $^{3}\text{He}$, we adopt the same constraint as in [7, 8]:

$$ (^{3}\text{He}/\text{D})_p < 0.59 \pm 0.27. $$

(A.2)

Here we have used the fact that the $^{3}\text{He}/\text{D}$ is a increasing function of the time since D is more destroyed than $^{3}\text{He}$ during chemical evolution, and hence the solar abundance gives the upper limit.

The $^4\text{He}$ is observed in metal poor extragalactic HII regions and the primordial abundance $Y_p$ is obtained by extrapolation of the observed abundances to zero metallicity. For a long
time relatively low abundance had been believed for $^4\text{He}$. Field and Olive [36] derived $Y_p = 0.238 \pm (0.002)_{\text{stat}} \pm (0.005)_{\text{syst}}$ and Izotov and Thuan [37] obtained $Y_p = 0.242 \pm 0.002$. However, recently, Olive and Skillman [38] reanalyzed the Izotov-Thuan data and derived higher $^4\text{He}$ abundance with much larger uncertainty, $Y_p = 0.249 \pm 0.09$. The similar value is obtained in more recent analysis [39] which we adopt here,

$$Y_p = 0.250 \pm 0.004.$$  \hspace{1cm} (A.3)

The smaller uncertainty than that in [38] is due to different determination of the electron temperature $T_e$ for HII regions. The OIII lines are used in [39] while [38] uses HeI recombination lines, which generally leads to large errors in $T_e$.

Warmest metal-poor (pop.II) halo stars have almost constant $^7\text{Li}$ abundances (Spite plateau) independent of metallicity. This constant value is considered as the primordial abundance of $^7\text{Li}$. Bonifacio et al. [40] obtained $\log_{10}(^7\text{Li}/\text{H})_p = -9.66 \pm 0.056$ which is consistent with more recent value $\log_{10}(^7\text{Li}/\text{H})_p = -9.63 \pm 0.06$ by Meléndez and Ramírez [41]. On the other hand, significantly small abundance $\log_{10}(^7\text{Li}/\text{H})_p = -9.91 \pm 0.10$ was derived in [42] where it was claimed that there is a correlation between $^7\text{Li}$ and Fe abundances due to $^7\text{Li}$ production by cosmic ray interactions and one should obtain the primordial value by taking zero Fe limit. The Fe dependence of $^7\text{Li}$ was also observed by Asplund et al [43] who reported $\log_{10}(^7\text{Li}/\text{H})_p = -9.90 \pm 0.06$. However, no correlation was observed in [40] [41]. Thus, the $^7\text{Li}$-Fe correlation is still an open question. Furthermore, the observed $^7\text{Li}$ abundance may be smaller than the primordial value if $^7\text{Li}$ is depleted in stars. For example, [41] shows that rotational mixing leads to depletion factor $D_7$ at most 0.3 dex. Therefore, at present it is difficult for us to reach some consensus. Since the upperbound of $^7\text{Li}$ is important in deriving a constraint on the sneutrino decay, we conservatively adopt the higher value [41] and add systematic error of 0.3 dex taking depletion into account,

$$\log_{10}(^7\text{Li}/\text{H})_p = -9.63 \pm 0.06 \pm 0.3.$$  \hspace{1cm} (A.4)

Here we have added $-0.3$ for systematic error because the high $^7\text{Li}$ abundances in [40] [41] are different from those in [42] [43] by about factor of 2 ($\simeq 0.3$ dex).

As for $^6\text{Li}$ recently Asplund et al. [43] detected $^6\text{Li}/^7\text{Li}$ ratios in 9 metal poor halo dwarfs. In particular, $^6\text{Li}$ abundance detected in very metal poor star LP 815-43 with $[\text{Fe/H}] = -2.74$ was is $^6\text{Li}/^7\text{Li} = 0.046 \pm 0.022$. Such high $^6\text{Li}$ abundance seems difficult to explain by the Galactic cosmic ray spallation and $\alpha$-fusion reactions and might be primordial. In this paper we regard it as upper bound on the primordial $^6\text{Li}$ abundance. Moreover, the depletion in stars is more important for $^6\text{Li}$ since $^6\text{Li}$ is more fragile than $^7\text{Li}$. The depletion factor $D_6$ for $^6\text{Li}$ is related to $D_7$ as $D_6 \simeq 2.5D_7$ [45], which leads to $\log_{10}(^6\text{Li}/^7\text{Li})_p = 1.5D_7 + \log_{10}(^6\text{Li}/^7\text{Li})_{\text{obs}}$. Taking account of the depletion, we adopt

$$(^6\text{Li}/^7\text{Li})_p < 0.046 \pm 0.022 + 0.084.$$  \hspace{1cm} (A.5)

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