Constraints on the Reheating Temperature in Gravitino Dark Matter Scenarios

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Considering gravitino dark matter scenarios, we study constraints on the reheating temperature of inflation. We present the gauge-invariant result for the thermally produced gravitino yield to leading order in the Standard Model gauge couplings. Within the framework of the constrained minimal supersymmetric Standard Model (CMSSM), we find a maximum reheating temperature of about $10^7$ GeV taking into account bound-state effects on the primordial $^6$Li abundance. We show that late-time entropy production can relax this constraint significantly. Only with a substantial entropy release after the decoupling of the lightest Standard Model superpartner, thermal leptogenesis remains a viable explanation of the cosmic baryon asymmetry within the CMSSM.

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INTRODUCTION

The observed flatness, isotropy, and homogeneity of the Universe suggest that its earliest moments were governed by inflation [1,2]. The inflationary expansion is followed by a phase in which the Universe is reheated. The reheating process repopulates the Universe and provides the initial conditions for the subsequent radiation-dominated epoch. We refer to the reheating temperature $T_R$ as the initial temperature of this early radiation-dominated epoch of our Universe.

The value of $T_R$ is an important prediction of inflation models. While we do not have evidence for temperatures of the Universe higher than $\mathcal{O}(1 \text{ MeV})$ (i.e., the temperature required by primordial nucleosynthesis), inflation models can point to $T_R$ well above $10^{10}$ GeV [2,3].

In this Letter we consider supersymmetric (SUSY) extensions of the Standard Model in which the gravitino $\tilde{G}$ is the lightest supersymmetric particle (LSP) and stable because of R-parity conservation. The gravitino LSP is a well-motivated dark matter candidate. As the gauge field of local SUSY transformations and the spin-3/2 superpartner of the graviton, the gravitino is an unavoidable implication of SUSY theories including gravity [3]. Its interactions are suppressed by inverse powers of the (reduced) Planck scale $M_P = 2.4 \times 10^{18}$ GeV. Its mass $m_{\tilde{G}}$ results from spontaneous SUSY breaking and can range from the eV scale up to scales beyond the TeV region [5].

While any initial population of gravitinos must be diluted away by the exponential expansion during inflation [6], gravitinos are regenerated in scattering processes of particles that are in thermal equilibrium with the hot primordial plasma. The efficiency of this thermal production of gravitinos during the radiation-dominated epoch is sensitive to $T_R$ [5,6,8,10,11]. Since the resulting gravitino density $\Omega_{\tilde{G}}^{\text{TP}}$ is bounded from above by the dark matter density $\Omega_{\text{dm}}$ [6], upper bounds on $T_R$ can be derived [5,12,13,14,15]. These bounds can be compared with predictions of the reheating temperature $T_R$ from inflation models. Moreover, $T_R$ is important for our understanding of the cosmic baryon asymmetry. For example, successful standard thermal leptogenesis [16] requires $T_R \gtrsim 10^9$ GeV [17].

We update the $T_R$ limits using the full gauge-invariant result for the relic density of thermally produced gravitinos, $\Omega_{\tilde{G}}^{\text{TP}}$, to leading order in the Standard Model gauge couplings [11]. This allows us to illustrate the dependence of the bounds on the gaugino-mass relation at the scale of grand unification $M_{\text{GUT}} \sim 2 \times 10^{16}$ GeV.

We consider gravitino dark matter scenarios also in the framework of the constrained minimal supersymmetric Standard Model (CMSSM) in which the gaugino masses, the scalar masses, and the trilinear scalar interactions are assumed to take on the respective universal values $m_{1/2}$, $m_0$, and $A_0$ at $M_{\text{GUT}}$. Taking into account gravitinos from thermal production and from late decays of the lightest Standard Model superpartner, we provide new upper bounds on the reheating temperature in the $(m_{1/2}, m_0)$ plane for various values of $m_{\tilde{G}}$. Previous studies of $T_R$ constraints within the CMSSM used the result of [10] to explore the viability of $T_R \gtrsim 10^9$ GeV [13,14]. Our study presents also scans for $T_R$ as low as $10^7$ GeV based on the result of [11] which includes electroweak contributions to thermal gravitino production [18].

In the considered CMSSM scenarios with the gravitino LSP, the next-to-lightest SUSY particle (NLSP) is either the lightest neutralino $\tilde{\chi}^0_1$ or the lighter stau $\tilde{\tau}_1$.1 Because of the extremely weak interactions of the gravitino, the NLSP typically has a long lifetime before it decays into the gravitino. If these decays occur during or after big-bang nucleosynthesis (BBN), the Standard Model particles emitted in addition to the gravitino can affect the abundance of the primordial light elements. Indeed, these BBN constraints disfavor the $\tilde{\chi}^0_1$ NLSP for $m_{\tilde{G}} \gtrsim 100$ MeV [13,14,21]. For the slepton NLSP case, the BBN constraints associated with

1 For simplicity, we consider $A_0 = 0$ in this work. For sizable $|A_0|$, also the lighter stop $\tilde{t}_1$ can be the NLSP [22,23].
hadronic/electromagnetic energy injection have also been estimated and found to be much weaker but still significant in much of the parameter space [13, 14, 15, 21].

Only recently, it has been stressed that bound–state formation of long-lived negatively charged particles with the primordial nuclei can affect BBN [22, 23, 24, 25]. With the charged long-lived stan NLS in, these bound–state effects also apply to the considered gravitino dark matter scenarios. In particular, a significant enhancement of \( \delta \) production after decoupling of the NLSP and before BBN can weaken the BBN constraints significantly [27]. At the same time, the gravitino density is estimated and found to be much weaker but still significant in much of the parameter space [13, 14, 15, 21].

Taking a conservative point of view, we do not include gravitino production in the radiation-dominated epoch, however, inflaton decays, for example, can lead to a sizable yield of non-thermally produced gravitinos depending on the inflation model; cf. [28, 29] and references therein.

**THERMAL GRAVITINO PRODUCTION**

Gravitinos with \( m_\tilde{G} \gtrsim 1 \text{ GeV} \) have decoupling temperatures of \( T_\text{f}^{\tilde{G}} \gtrsim 10^{14} \text{ GeV} \), as will be shown below. We consider thermal gravitino production in the radiation-dominated epoch starting at \( T_R < T_\text{f}^{\tilde{G}} \), assuming that inflation has diluted away any initial gravitino population.\(^2\) For \( T_R < T_\text{f}^{\tilde{G}} \), gravitinos are not in thermal equilibrium with the post-inflatonium plasma. Accordingly, the evolution of the gravitino number density \( n_\tilde{G} \), with cosmic time \( t \) is described by the following Boltzmann equation [11]

\[
\frac{dn_\tilde{G}}{dt} + 3H n_\tilde{G} = C_\tilde{G}
\]  

\[
C_\tilde{G} = \sum_{i=1}^{3} \frac{3\zeta_i^T T^6}{16\pi^2 M_P^2} \left( 1 + \frac{M^2}{3m_\tilde{G}^2} \right) c_i g_i^2 \ln \left( \frac{k_i}{g_i} \right) ,
\]

where \( H \) denotes the Hubble parameter. The collision term \( C_\tilde{G} \) involves the gaugino mass parameters \( M_i \), the gaugino couplings \( g_i \), and the constants \( c_i \) and \( k_i \), associated with the gauge groups \( U(1)_Y, SU(2)_L \), and \( SU(3)_c \) as given in Table I. In expression (2) the temperature \( T \) provides the scale for the evaluation of \( M_i \) and \( g_i \). The given collision term is valid for temperatures sufficiently below the gravitino decoupling temperature, where gravitino disappearance processes can be neglected. A primordial plasma with the particle content of the minimal SUSY Standard Model (MSSM) in the high-temperature limit is used in the derivation of (2).

The collision term (2) results from a consistent gauge-invariant finite-temperature calculation [11, 18] following the approach used in Ref. [10]. Thus, in contrast to the previous estimates in [7, 8], the expression for \( C_\tilde{G} \) is independent of arbitrary cutoffs. Note that the field-theoretical methods of [30, 31] applied in its derivation require weak couplings, \( g_i \ll 1 \), and thus high temperatures \( T \gg 10^6 \text{ GeV} \).

Assuming conservation of entropy per comoving volume, the Boltzmann equation (1) can be solved to good approximation analytically [10, 32]. At a temperature \( T_\text{low} \ll T_R \), the resulting gravitino yield from thermal production reads

\[
Y_\text{TP}^{\tilde{G}}(T_\text{low}) \equiv \frac{n_\tilde{G}^{\text{TP}}(T_\text{low})}{s(T_\text{low})} \approx \frac{C_\tilde{G}(T_R)}{s(T_R) H(T_R)}
\]

\[
= \sum_{i=1}^{3} y_i g_i^2(T_R) \left( 1 + \frac{M_i^2(T_R)}{3m_\tilde{G}^2} \right) \times \ln \left( \frac{k_i}{g_i(T_R)} \right) \left( \frac{T_R}{10^{10}\text{GeV}} \right)^{3} ,
\]

where the constants \( y_i \) are given in Table I. These constants are obtained with the Hubble parameter describing the radiation-dominated epoch, \( H_{\text{rad}}(T) =
the entropy density \( s(T) = \sqrt{g_\ast(T)} \pi^2/90 T^2/\mathrm{M}_\mathrm{P} \), the entropy density \( s(T) = 2 \pi^2 g_\ast s(T) T^3/45 \), and an effective number of relativistic degrees of freedom of \( g_\ast(T_R) = g_\ast s(T_R) = 228.75 \). We evaluate \( g_\ast(T_R) \) and \( M_\ast(T_R) \) using the one-loop evolution described by the renormalization group equation in the MSSM:

\[
g_\ast(T) = \left( g_\ast^2(m_{\tilde{Z}}) - \frac{\beta_2(1)}{8 \pi^2} \ln \left( \frac{T}{m_{\tilde{Z}}} \right) \right)^{-1/2}, \quad (4)
\]

\[
M_\ast(T) = \left( \frac{g_\ast(T)}{g_\ast(M_{\mathrm{GUT}})} \right)^2 M_\ast(M_{\mathrm{GUT}}) \quad (5)
\]

with the respective gauge coupling at the Z-boson mass, \( g_\ast(m_{\tilde{Z}}) \), and the \( \beta_2(1) \) coefficients listed in Table I.

Without late-time entropy production, the gravitino yield from thermal production at the present temperature \( T_0 \) is given by

\[
Y^{\text{TP}}(T_0) = Y^{\text{TP}}(T_{\text{low}}). \quad (6)
\]

The resulting density parameter of thermally produced gravitinos is

\[
\Omega^{\text{TP}} h^2 = m_{\tilde{G}} Y^{\text{TP}}(T_0) \frac{s(T_0) h^2}{\rho_c} \quad (7)
\]

with the Hubble constant \( h \) in units of 100 km Mpc\(^{-1}\) s\(^{-1}\) and \( \rho_c/\rho(s(T_0) h^2) = 3.6 \times 10^{-9} \) GeV.

In Fig. 1 our result \( 5 \) for the thermally produced gravitino yield \( Y^{\text{TP}}(T_{\text{low}}) \) is shown as a function of \( T_R \) for various values of \( m_{\tilde{G}} \) (solid lines). The curves are obtained with \( m_{1/2} = 500 \) GeV for the case of universal gaugino masses at \( M_{\mathrm{GUT}} \): \( M_{1,2,3}(M_{\mathrm{GUT}}) = m_{1/2} \). The dotted lines show the corresponding results from the SU(3)c yield of Ref. [10] for \( M_3 = m_{1/2} \), which was used to study \( T_R \) constraints on gravitino dark matter scenarios in Refs. [13, 14, 15]. We find that \( 5 \) exceeds the yield derived from [10] by about 50%; cf. [11]. The dashed (blue in the web version) horizontal line indicates the equilibrium yield

\[
Y^{\text{eq}}(T_R) \equiv \frac{n_{\tilde{G}}^{\text{eq}}}{s} \approx 1.8 \times 10^{-3} \quad (8)
\]

which is given by the equilibrium number density of a relativistic spin 1/2 Majorana fermion, \( n_{\tilde{G}}^{\text{eq}} = 3 \zeta(3)/2 \pi^2 \). For \( T > T_{\tilde{G}}^0 \), \( g_\ast(T) = g_\ast s(T) = 230.75 \) since the spin 1/2 components of the gravitinos are in thermal equilibrium. In the region where the yield \( 5 \) approaches the equilibrium value \( 5 \), gravitino disappearance processes should be taken into account. This would then lead to a smooth approach of the non-equilibrium yield to the equilibrium abundance. Without the back-reactions taken into account, the kink position indicates a lower bound for \( T_{\tilde{G}}^0 \). Towards smaller \( m_{\tilde{G}} \), \( T_{\tilde{G}}^0 \) decreases due to the increasing strength of the gravitino couplings. For example, for \( m_{\tilde{G}} = 1 \) GeV (10 MeV), we find \( T_{\tilde{G}}^0 \gtrsim 10^{14} \) GeV (10\(^{10} \) GeV).

In the analytical expression \( 5 \) we refer to \( T_R \) as the initial temperature of the radiation-dominated epoch. So far we have not considered the phase in which the coherent oscillations of the inflaton field \( \phi \) dominate the energy density of the Universe, where one usually defines \( T_R \) in terms of the decay width \( \Gamma_\phi \) of the inflaton field \( \phi \). To account for the reheating phase, we numerically integrate \( 5 \) together with the Boltzmann equations for the energy densities of radiation and the inflaton field,

\[
\frac{dp_{\text{rad}}}{dt} + 4H p_{\text{rad}} = \Gamma_\phi \rho_\phi, \quad (9)
\]

\[
\frac{dp_\phi}{dt} + 3H p_\phi = -\Gamma_\phi \rho_\phi, \quad (10)
\]

respectively; for details see Appendix F of Ref. [34].

With our result for the collision term \( 2 \), we find that the gravitino yield obtained numerically is in good agreement with the analytical expression \( 5 \) for

\[
T_R \simeq \left[ \frac{90}{g_\ast(T_R)^2} \right]^{1/4} \sqrt{\frac{\Gamma_\phi M_\mathrm{P}}{1.8}} \quad (11)
\]

which satisfies \( \Gamma_\phi \gtrsim 1.8 H_{\text{rad}}(T_R) \). For an alternative \( T_R \) definition given by \( \Gamma_\phi = \xi H_{\text{rad}}(T_R) \),

\[
T_{\xi}^R \equiv \left[ \frac{90}{g_\ast(T_R)^2} \right]^{1/4} \sqrt{\frac{\Gamma_\phi M_\mathrm{P}}{\xi}}, \quad (12)
\]

![FIG. 1: The thermally produced gravitino yield (5) as a function of \( T_R \) for \( m_{\tilde{G}} = 10 \) MeV, 100 MeV, 1 GeV, 10 GeV, 100 GeV, and 1 TeV (solid lines from left to right) and \( M_{1,2,3}(M_{\mathrm{GUT}}) = m_{1/2} = 500 \) GeV. The dotted lines show the corresponding yield obtained with the SU(3)c result for the collision term of Ref. [10]. The dashed (blue in the web version) horizontal line indicates the equilibrium yield of a relativistic spin 1/2 Majorana fermion.](image)
the associated numerically obtained gravitino yield is described by the analytical expression obtained after substituting $T_R$ with $\sqrt{\xi/1.8} T_R$ in [39].

While we focus on scenarios in which the gravitino is stable, the yield [33] is also crucial to extract cosmological constraints in scenarios with unstable gravitinos. Based on the result of [10] and taking into account thermal gravitino production during reheating, the following fitting formula was used to study constraints from decaying gravitinos in Refs. [33, 34, 35]:

$$Y_{\tilde{G}}^{KRM}(T_{\text{low}}) \simeq 1.9 \times 10^{-12} \left( \frac{T_R}{10^{10} \text{ GeV}} \right)^3 \left[ 1 + 0.045 \ln \left( \frac{T_R}{10^{10} \text{ GeV}} \right) \right] \times \left[ 1 - 0.028 \ln \left( \frac{T_R}{10^{10} \text{ GeV}} \right) \right], \quad (13)$$

where $T_R$ was defined via $\Gamma_\phi = 3 H_{\text{rad}}(T_R)$. Comparing [13] with our result after the matching of the $T_R$ definitions, we find that our result exceeds the $m_{\tilde{G}}$-independent yield [13] by about 30% for $m_{\tilde{G}} \gg M_i(T_R)$. While the $m_{\tilde{G}}$ dependence of $Y_{\tilde{G}}^{\text{TP}}$ becomes negligible for decreasing $M_i(T_R)/m_{\tilde{G}}$, the yield [13] is used for $m_{\tilde{G}}$ as small as 100 GeV in Refs. [33, 34, 35]. As can be seen in Fig. 1, the actual yield for $m_{\tilde{G}} = 100$ GeV is thereby underestimated by about an order of magnitude. Accordingly, the $T_R$ bounds given in [33, 34, 35] are underestimated in the region $m_{\tilde{G}} < 1$ TeV.

CONSTRANTS ON $T_R$

The reheating temperature $T_R$ is limited from above in the case of a stable gravitino LSP since $\Omega_{\tilde{G}}^{\text{TP}}$ cannot exceed the dark matter density $\Omega_{\text{dm}}^G$ [8, 12, 13, 14, 15]. In this paper, we use [36, 37]:

$$\Omega_{\text{dm}}^G h^2 = 0.105^{+0.021}_{-0.030} \quad (14)$$

as obtained from the measurements of the cosmic microwave background (CMB) anisotropies by the Wilkinson Microwave Anisotropy Probe (WMAP) satellite.\(^3\)

In Fig. 2 we show the resulting upper limits on $T_R$ as a function of $m_{\tilde{G}}$. On the gray band, the thermally produced gravitino density [7] is within the nominal 3σ range [14]. The upper (lower) gray band is obtained for $M_{1,2,3} = m_{1/2}$ at $M_{\text{GUT}}$ with $m_{1/2} = 500$ GeV (2 TeV). The dashed lines show the corresponding constraints for the exemplary non-universal scenario [39].

$M_{1/2} = 500$ GeV in Fig. 2 is by about a factor of 4 more severe than the $T_R$ limits shown in Figs. 5 and 6 of Ref. [14] in the region $m_{\tilde{G}} \lesssim 10$ GeV in which $\Omega_{\tilde{G}}^{\text{TP}}$ governs the limits. It seems to us that the gluino mass in (1.2) of Ref. [14] was accidentally evaluated at the scale $\mu = T_R$ rather than at the scale $\mu \simeq 100$ GeV; see Sec. 5 of [10].

FIG. 2: Upper limits on the reheating temperature $T_R$. On the upper (lower) gray band, $\Omega_{\tilde{G}}^{\text{TP}}$ for $M_{1,2,3} = m_{1/2} = 500$ GeV (2 TeV) at $M_{\text{GUT}}$ agrees with $\Omega_{\text{dm}}^G$. The corresponding $T_R$ limits from the requirement $\Omega_{\tilde{G}}^{\text{TP}} h^2 \leq 0.126$ shown by the dashed and dotted lines are obtained respectively with [39] for $M_{1/2} = 500$ GeV at $M_{\text{GUT}}$ and with the result of Ref. [10] for $M_{3} = m_{1/2}$ at $M_{\text{GUT}}$.

\(^3\) This nominal 3σ range is derived assuming a restrictive six-parameter “vanilla” model. A larger range is possible—even with additional data from other cosmological probes—if the fit is performed in the context of a more general model [35].
CONSTRAINTS ON $T_R$ IN THE CMSSM

In the CMSSM, one assumes universal soft SUSY breaking parameters at $M_{GUT}$. The CMSSM yields phenomenologically acceptable spectra with only four parameters and a sign: the gaugino mass parameter $m_{1/2}$, the scalar mass parameter $m_0$, the trilinear coupling $A_0$, the mixing angle $\tan \beta$ in the Higgs sector, and the sign of the higgsino mass parameter $\mu$.

Assuming $A_0 = 0$ for simplicity, the lightest Standard Model superpartner is either the lightest neutralino $\tilde{\chi}_1^0$ or the lighter stau $\tilde{\tau}_1$. Indeed, most CMSSM investigations assume that $\tilde{\chi}_1^0$ is the LSP that provides dark matter; cf. [10] and references therein. The parameter region in which $m_{\tilde{\tau}_1} < m_{\tilde{\chi}_1^0}$ is usually not considered because of the severe upper limits on the abundance of stable charged particles [37]. In the gravitino LSP case, $m_{\tilde{\tau}_1} < m_{\tilde{\chi}_1^0}$ can be viable since the lightest Standard Model superpartner is unstable [13, 14, 25, 11] and references therein.

With the gravitino LSP, the lightest Standard Model superpartner is the NLSP that decays into Standard Model particles and one gravitino LSP. For $m_{\tilde{G}} \geq 1$ GeV, the NLSP decays after its decoupling from the thermal plasma. Thus, the relic density of the associated nonthermally produced gravitinos reads [42]

$$\Omega_{\tilde{G}}^\text{NTP} h^2 = \frac{m_{\tilde{G}} Y_{\text{NLSP}}(T_0) s(T_0) h^2/\rho_c}{m_{\text{NLSP}}} = \frac{m_{\tilde{G}}}{m_{\text{NLSP}}} \Omega_{\text{NLSP}} h^2,$$

where $m_{\text{NLSP}}$ is the mass of the NLSP and $Y_{\text{NLSP}}(T_0)$ and $\Omega_{\text{NLSP}} h^2$ are respectively the yield and the relic density that the NLSP would have today, if it had not decayed.

In Fig. 3, the solid (black) and dotted (blue in the web version) lines show respectively contours of $Y_{\text{NLSP}}(T_0)$ and $\Omega_{\text{NLSP}} h^2$ are respectively the yield and the relic density that the NLSP would have today, if it had not decayed. The medium gray and the light gray regions at small $m_{1/2}$ are excluded respectively by the mass bounds $m_{\tilde{\chi}_1^0} > 94$ GeV and $m_{11/2} > 114.4$ GeV from chargino and Higgs searches at LEP [37]. The leftmost dotted line (blue in the web version) line indicates the LEP bound $m_{\tilde{\tau}_1} > 81.9$ GeV [37]. For $\tan \beta = 30$, tachyonic sfermions occur in the low-energy spectrum at points in the white region labeled as “tachyonic.” We employ the FORTRAN program SuSpect [13] to calculate the low-energy spectrum of the superparticles and the Higgs bosons, where we use $m_t = 172.5$ GeV for the top quark mass. Assuming standard cosmology, the yield $Y_{\text{NLSP}}(T_0)$ is obtained from the $\Omega_{\text{NLSP}} h^2$ values provided by the computer program micrOMEGAs [11].

The contours shown in Fig. 3 are independent of $m_{\tilde{G}}$ and $T_R$. Therefore, they can be used to interpret the results shown in the figures below. Note the sensitivity of both $Y_{\tilde{\chi}_1^0}(T_0)$ and $m_{\tilde{\tau}_1}$ on $\tan \beta$. By going from $\tan \beta = 10$ to $\tan \beta = 30$, $Y_{\tilde{\chi}_1^0}(T_0)$ decreases by about a factor of two at points that are not in the vicinity of the dashed line, i.e., that are outside of the $\tilde{\tau}_1-\tilde{\chi}_1^0$ coannihilation region. While $m_{\tilde{\tau}_1}$ becomes smaller by increasing $\tan \beta$ to 30, the $\tan \beta$ dependence of $m_{\tilde{\chi}_1^0}$ is negligible.

Let us now explore the parameter space in which

$$0.075 \leq \Omega_{\tilde{G}}^\text{NTP} h^2 + \Omega_{\text{NLSP}}^\text{NTP} h^2 \leq 0.126 .$$

Now, $T_R$ and $m_{\tilde{G}}$ appear in addition to the traditional CMSSM parameters. We focus on $m_{\tilde{G}} \geq 1$ GeV since the soft SUSY breaking parameters of the CMSSM are usually assumed to result from gravity-mediated SUSY breaking. However, we do not restrict our study to fixed relations between $m_{\tilde{G}}$ and the soft SUSY breaking parameters such as the ones suggested, for example, by the Polonyi model.

In Fig. 4, the light, medium, and dark shaded (green in the web version) bands show the $(m_{1/2}, m_0)$ regions that satisfy the constraint (17) for $T_R = 10^7$, $10^8$, and $10^9$ GeV, respectively, where $\tan \beta = 10$, $A_0 = 0$, $\mu > 0$. The four panels are obtained for the choices (a) $m_{\tilde{G}} = 10$ GeV, (b) $m_{\tilde{G}} = 100$ GeV, (c) $m_{\tilde{G}} = 0.2 m_0$, and (d) $m_{\tilde{G}} = m_0$. In the dark-gray region, the gravitino is not the LSP. The regions excluded by the chargino and Higgs mass bounds and the line indicating $m_{\tilde{\chi}_1^0} = m_{\tilde{\tau}_1}$ are identical to the ones shown in the left panel of Fig. 3. The dotted lines show contours of the NLSP lifetime. For the $\tilde{\tau}_1$ NLSP,

$$\tau_{\tilde{\tau}_1} \simeq \Gamma^{-1}(\tilde{\tau}_1 \to \tilde{G}\tau) = \frac{48\pi m_{\tilde{\tau}_1}^2 M_\text{Pl}^2}{m_{\tilde{\tau}_1}^2} \left(1 - \frac{m_{\tilde{\tau}_1}^2}{m_{\tilde{G}}^2}ight)^{-4}$$

as obtained in the limit $m_{\tilde{\tau}_1} \to 0$. For the $\tilde{\chi}_1^0$ NLSP, we calculate $\tau_{\tilde{\chi}_1^0}$ from the expressions given in Sec. IIC of Ref. [21].

The $\tau_{\text{NLSP}}$ contours in Fig. 4 illustrate that the NLSP decays during/after BBN. Successful BBN predictions therefore imply cosmological constraints on $m_{\tilde{G}}$, $m_{\text{NLSP}}$, and $Y_{\text{NLSP}}$ [13, 14, 15, 21]. Indeed, it has been found that the considered $\tilde{\chi}_1^0$ NLSP region is completely disfavored for $m_{\tilde{G}} \geq 1$ GeV by constraints from late electromagnetic and hadronic energy injection [13, 14, 21, 22]. In the $\tilde{\tau}_1$ NLSP region, the constraints from electromagnetic and hadronic energy release are important but far less severe than in the $\tilde{\chi}_1^0$ NLSP case. Thus, much of the $\tilde{\tau}_1$ NLSP region was believed to be cosmologically allowed [13, 14, 15, 21, 22].

Recently, this picture has changed. It has been found that bound-state formation of long-lived negatively

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5 The NLSP freezeout temperature can be estimated from its mass: $T^\text{NLSP} \lesssim m_{\text{NLSP}}/20$. Thus, $T_R \gg T^\text{NLSP}$ for $T_R > 10^9$ GeV which is considered in this Letter.
charged $\tilde{\tau}_1$'s with primordial nuclei can catalyze the production of $^6$Li significantly\cite{22,25}. Indeed, in most of the $\tilde{\tau}_1$ NLSP parameter space, the associated bounds are much more severe than the ones from late energy injection. Only for $\tau_{\tilde{\tau}_1} \lesssim 10^3$ s and $m_{\tilde{\tau}_1} \gtrsim 40$ GeV, the constraints from hadronic energy release can become more severe than the ones from catalyzed $^6$Li production\cite{25,20}. We thus consider both the constraint from catalyzed $^6$Li production derived in\cite{22} and the one from late hadronic energy injection derived in\cite{15,9}.

For the constraint from bound-state effects on $^6$Li production, we adopt the bounds given in Fig. 4 of Ref.\cite{22} as $\tau_{\tilde{\tau}_1}$-dependent upper limits on the yield of the negatively charged staus, $Y_{\text{NLSP}}/2$. These bounds are obtained assuming a limiting primordial abundance of \cite{13}

$$\left(\frac{^6\text{Li}}{^4\text{He}}\right)_p \lesssim 2 \times 10^{-11}. \quad (19)$$

The resulting constraint disfavors the $\tilde{\tau}_1$ NLSP region to the left of the long-dash-dotted (red in the web version) line shown in Fig.\ref{fig:3}

For the constraint from late hadronic energy injection, we use the upper limits on $Y_{\text{NLSP}}$ that are given in Fig. 11 of Ref.\cite{17}. These limits are derived from a computation of the 4-body decay of the stau NLSP into the gravitino, the tau, and a quark-antiquark pair.\footnote{For details on the other BBN bounds and the additional CMB bounds, we refer the reader to the detailed investigations presented in Refs.\cite{12,13,21,41,46}. They are based on the severe and conservative upper bounds on the released hadronic energy (95% CL) obtained in\cite{34} for observed values of the primordial D abundance of

$$\left(\frac{n_D}{n_H}\right)_{\text{mean}} = (2.78^{+0.44}_{-0.26}) \times 10^{-5} \quad (\text{severe}),$$
$$\left(\frac{n_D}{n_H}\right)_{\text{high}} = (3.98^{+0.59}_{-0.67}) \times 10^{-5} \quad (\text{conservative}).$$

In Fig.\ref{fig:3} the associated constraints are shown by the short-dash-dotted (blue in the web version) lines. The D constraint disfavors the region between the corresponding lines in panel (b) and the region above the corresponding lines in panels (c) and (d). In panel (a) the D constraint does not appear.\footnote{The 3-body estimate of the hadronic energy release given in Ref.\cite{22} leads to overly restrictive limits, as shown in Ref.\cite{13}. Additional constraints on hadronic energy release are imposed by the primordial abundances of $^4$He, $^3$He/$^4$He, $^7$Li, and $^4$Li/$^7$Li\cite{21,34,45,48,49}. However, in the region allowed by the $^4$Li constraint from bound-state effects, i.e., $\tau_{\tilde{\tau}_1} \lesssim 10^3$ s, the considered D constraint on hadronic energy release is the dominant one as can be seen in Figs. 38–41 of Ref.\cite{34} and Figs. 6–8 of Ref.\cite{50}.}

Remarkably, one finds in each panel of Fig.\ref{fig:3} that the highest $T_R$ value allowed by the considered BBN constraints is about $10^7$ GeV. The bands obtained for $T_R \gtrsim 10^8$ GeV are located completely within the region...}

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig3.png}
\caption{Contours of $Y_{\text{NLSP}}(T_R)$ (solid black lines) and $m_{\text{NLSP}}$ (dotted blue lines) in the $(m_{1/2}, m_0)$ plane for $A_0 = 0$, $\mu > 0$, $\tan \beta = 10$ (left panel) and $\tan \beta = 30$ (right panel). Above (below) the dashed line, $m_{\tilde{\tau}_1} < m_{\tilde{\chi}_1^0}$ ($m_{\tilde{\tau}_1} < m_{\tilde{\chi}_1^0}$). The medium gray and the light gray regions at small $m_{1/2}$ show the mass bounds $m_{\tilde{\chi}_1^\pm} > 94$ GeV and $m_H > 114.4$ GeV from chargino and Higgs searches at LEP\cite{37}.}
\label{fig:3}
\end{figure}
$m_\chi = 10$ GeV, $m_{1/2} = 100$ GeV, and $m_{Higgs} = m_0$. In each panel, the light, medium, and dark shaded (green in the web version) bands indicate the regions in which $0 < \Omega h^2 \leq 0.126$ for $T_e = 10^7$, $10^8$, and $10^9$ GeV, respectively. The medium gray and the light gray regions at small $m_{1/2}$ are excluded respectively by chargino and Higgs searches at LEP. In the dark gray region, the gravitino is not the LSP. The dotted lines show contours of the NLSP lifetime. Below the dashed line, $m_{\chi_1} < m_{\chi_2}$. With the $\bar{\chi}_1$ NLSP, the region to the left of the long-dash-dotted (red in the web version) line is cosmologically disfavored by bound-state effects on the primordial $^6$Li abundance. The effects of late hadronic energy injection on the primordial D abundance disfavor the $\bar{\chi}_1$ NLSP region between the short-dash-dotted (blue in the web version) lines in panel (b) and the one above the corresponding lines in panels (c) and (d). The $\chi_1^0$ NLSP region above the dashed line, in which $m_{\chi_2} < m_{\chi_1}$, is cosmologically disfavored by the effects of late electromagnetic/hadronic energy injection on the abundances of the light primordial elements.
disfavored by the $^6$Li bound. In previous gravitino dark matter studies within the CMSSM that did not take into account bound-state effects on the primordial $^6$Li abundance, much higher temperatures of up to about $10^9$ GeV were believed to be allowed [11, 13, 14].

The constraint $T_R \lesssim 10^7$ GeV remains if we consider larger values of $\tan \beta$. This is demonstrated in Fig. 5 for $\tan \beta = 30$, $A_0 = 0$, $\mu > 0$, and $m_\chi = m_0$. While the $^6$Li constraint relaxed in this way would appear in Fig. 4(b) as an almost vertical line slightly above $m_{1/2} = 3$ TeV.

While the constraint $T_R \lesssim 10^7$ GeV is found for each of the considered $m_\chi$ relations, one cannot use the $^6$Li bound to set bounds on $m_\tau_1$ without insights into $m_\chi$. The $^6$Li bound disappears for $\tau_1 \lesssim 10^3 \text{ s}$ [22] which is possible even for $m_\tau_1 = O(100 \text{ GeV})$ provided $m_\chi$ is sufficiently small; see [15]. However, the constraints on $T_R$ become more severe towards small $m_\chi$ as is shown in Fig. 4. Thus, the constraint $T_R \lesssim 10^7$ GeV cannot be evaded by lowering $m_\chi$ provided $T_R < T_f^G$.

An upper limit on $T_R$ of $10^7$ GeV can be problematic for inflation models and baryogenesis scenarios. This finding can thus be important for our understanding of the thermal history of the Universe.

**CONSTRAINTS ON $T_R$ WITH LATE-TIME ENTROPY PRODUCTION**

The constraints shown above are applicable for a standard thermal history during the radiation-dominated epoch. However, it is possible that a substantial amount of entropy is released, for example, in out-of-equilibrium decays of a long-lived massive particle species $X$ [2, 51].

If $X$ lives sufficiently long, it might decay while its rest mass dominates the energy density of the Universe. The associated evolution of the entropy per comoving volume, $S \equiv s a^3$, is described by [2, 51]

$$\frac{dS}{dt} = \frac{\Gamma_X \rho_X a^3}{T} = \left(\frac{2 \pi^2}{45} g_* \right)^{1/3} \Gamma_X \rho_X a^4 S^{-1/3}$$

(20)

together with the Boltzmann equation (10) for $\phi = X$ and the Friedmann equation governing the evolution of the scale factor of the Universe $a$. Here $\Gamma_X$ and $\rho_X$ denote respectively the decay width and the energy density of $X$. Thus, the temperature after the decay can be expressed in terms of $\Gamma_X$,

$$T_{\text{after}} \equiv \left[ \frac{10}{g_* (T_{\text{after}})^{\frac{1}{2}}} \right]^{1/4} \sqrt{\Gamma_X M_P} ,$$

(21)

which satisfies $\Gamma_X = 3 H_{\text{rad}} (T_{\text{after}})$. Indeed, primordial nucleosynthesis imposes a lower limit on this temperature $[52, 53, 54, 55]$

$$T_{\text{after}} \gtrsim 0.7 \sim 4 \text{ MeV}.$$  

(22)

In Fig. 4 we show the evolution of $S$, $a^3 \rho_X$, and $a^3 \rho_{\text{rad}}$ for two exemplary scenarios respecting (22). The scale factor $a$ is normalized by $a_0 \equiv a(10 \text{ GeV}) = 1 \text{ GeV}^{-1}$ and the temperature dependence of $g_*$ is taken into account as determined in [53]. For $\rho_X(10 \text{ GeV}) = 0.1 \rho_{\text{rad}}(10 \text{ GeV})$ and $T_{\text{after}} = 6 \text{ MeV}$, $S$ increases by a factor of $\Lambda = 100$ as shown by the corresponding solid line. For $\rho_X(10 \text{ GeV}) = 8 \rho_{\text{rad}}(10 \text{ GeV})$ and $T_{\text{after}} = 4.9 \text{ MeV}$, $S$ increases by a factor of $\Delta = 10^4$ as shown by the corresponding dotted (blue in the web version) line.

---

9 Note that our bands for $T_R = 10^9$ GeV differ from the ones shown in Refs. [12] [13]; see footnote 4.
The constraints discussed below shall therefore be considered as conservative bounds. For studies of gravitino production during an entropy producing event, we refer to Fig. 6 and references therein.

Conversely, in the case of late-time entropy production after the decoupling of the NLSP (and before BBN both, $Y^{\text{TP}}_G(T_0)$ and $Y_{\text{NLSP}}(T_0)$, are reduced:

$$Y^{\text{TP}}_G(T_0) = \frac{1}{\Delta} Y^{\text{TP}}_G(T_{\text{low}})$$

$$Y_{\text{NLSP}}(T_0) = \frac{1}{\Delta} Y_{\text{NLSP}}(T_{\text{low}}) \quad (25)$$

Accordingly, $\Omega^{\text{TP}}_G$ and $\Omega^{\text{NTP}}_{G,N}$ become smaller and the BBN constraints can be relaxed.

In Fig. 7 we show how late-time entropy production before (left) and after (right) NLSP decoupling affects the $^{6}\text{Li}$ bound in the somewhat narrow window between NLSP decoupling and BBN. This is different for the viability of thermal leptogenesis in the considered scenarios ($T_{\tilde{f}} > T_R$) and for collider prospects as discussed below.

In the case of late-time entropy production before the decoupling of the NLSP, we parameterize this by writing

$$Y^{\text{TP}}_G(T_0) = \frac{1}{\delta} Y^{\text{TP}}_G(T_{\text{low}}) \quad (24)$$

In this case, $Y_{\text{NLSP}}(T_0)$ and thereby $\Omega^{\text{NTP}}_G$ and the BBN constraints remain unaffected.

We restrict our study to entropy production at late times, $T_{\text{before}} \approx T_{\text{low}} \ll T_R$, so that the thermal production of gravitinos is not affected. To work in a model independent way, we assume that the production of gravitinos and NLSPs in the entropy producing event, such as the direct production in decays of $X$, is negligible. Moreover, in this section, we focus on scenarios in which the decoupling of the NLSP is not or at most marginally affected by entropy production, i.e., either $T_R \gg T_{\text{after}} \gg T_{\text{NLSP}}^\text{low}$ or $\rho_{\text{rad}} \gg \rho_X$ for $T \gtrsim T_{\text{NLSP}}^\text{low}$. Thus, the thermally produced gravitino yield and—in the case of entropy production after NLSP decoupling—also the non-thermally produced gravitino yield are diluted:

$$Y^{\text{TP}}_G(T_{\text{after}}) = \frac{S(T_{\text{low}})}{S(T_{\text{after}})} Y^{\text{TP}}_G(T_{\text{low}}) \quad (23)$$

In the case of late-time entropy production before the decoupling of the NLSP, we parameterize this by writing

$$Y^{\text{TP}}_G(T_{\text{after}}) = \frac{S(T_{\text{low}})}{S(T_{\text{after}})} Y^{\text{TP}}_G(T_{\text{low}}) \quad (24)$$

In this case, $Y_{\text{NLSP}}(T_0)$ and thereby $\Omega^{\text{NTP}}_G$ and the BBN constraints remain unaffected.

11 The constraints discussed below shall therefore be considered as conservative bounds. For studies of gravitino production during an entropy producing event, we refer to Fig. 6 and references therein.

FIG. 6: Evolution of $S$, $a^2\rho_X$, and $a^2\rho_{\text{rad}}$ as a function of $T$ for the normalization $a \equiv a(10 \text{ GeV}) = 1 \text{ GeV}^{-1}$. The solid lines are obtained for $\rho_X(10 \text{ GeV}) = 0.1 \rho_{\text{rad}}(10 \text{ GeV})$ and $T_{\text{after}} = 6 \text{ MeV}$, the dotted (blue in the web version) lines for $\rho_X(10 \text{ GeV}) = 8 \rho_{\text{rad}}(10 \text{ GeV})$ and $T_{\text{after}} = 4.9 \text{ MeV}$.

Note that we do not show the D constraint on $^{6}\text{Li}$ bound is independent of $\delta$ and vanishes already for $\Delta = 10$; an exception is the severe D constraint which still appears for $\Delta = 10$ in panel (d). BBN constraints on $\Omega^{\text{NTP}}_G$ NLSP scenarios with entropy production after NLSP decoupling will be studied elsewhere.

Comparing panels (b) and (d) of Fig. 4 with panels (a) and (c) in Fig. 7, we find that a dilution factor of $\delta = 10$ (100) relaxes the $T_R$ bound by a factor of 10 (100). Since the BBN constraints are unaffected by $\delta$, the cosmologically disfavored range of NLSP masses cannot be relaxed. With the dilution after NLSP decoupling, the relaxation of the $T_R$ constraints is more pronounced. Here also the cosmologically disfavored range of NLSP masses can be relaxed. However, as can be seen in panels (b) and (d) of Fig. 7, the $^{6}\text{Li}$ bound is persistent. With a dilution factor of $\delta = 100$, large regions of the $(m_{\tilde{\chi}/m_0})$ plane remain cosmologically disfavored. For $\Delta \gtrsim 10^4$, however, the $^{6}\text{Li}$ bound can be evaded as will be shown explicitly below.

Figure 7 shows that inflation models predicting, for example, $T_R = 10^9 \text{ GeV}$ become allowed in the CMSSM with gravitino dark matter for $\delta = \Delta \approx 100$. Here it is not necessary to have late-time entropy production in the somewhat narrow window between NLSP decoupling and BBN. This is different for the viability of thermal leptogenesis in the considered scenarios ($T_{\tilde{f}} > T_R$) and for collider prospects as discussed below.
FIG. 7: The effect of late-time entropy production before (left) and after (right) NLSP decoupling on regions in which $0.075 \leq \Omega_S h^2 \leq 0.126$ for $T_R = 10^9$ GeV. The $(m_{1/2}, m_0)$ plane is shown for $\tan \beta = 10, A_0 = 0, \mu > 0, m_{\tilde{G}} = 100$ GeV (upper panels) and $m_{\tilde{G}} = m_0$ (lower panels). The dark shaded (green in the web version) region is obtained without late-time entropy production $\delta = \Delta = 1$. The medium and light shaded (green in the web version) bands are obtained with a dilution of $\Omega_G^T$ ($\Omega_G^T + \Omega_G^{\text{NTTP}}$) by $\delta = 10$ ($\Delta = 10$) and $\delta = 100$ ($\Delta = 100$), respectively. The $\tilde{\tau}_1$ NLSP region to the right of the dot-dashed (red in the web version) line is cosmologically disfavored by the primordial $^6$Li abundance. Other curves and regions are identical to the ones in the corresponding panels of Fig. 4. The severe D constraint for $\Delta = 10$ appears only in panel (d).

THERMAL LEPTOGENESIS IN THE CMSSM WITH GRAVITINO DARK MATTER

The constraint $T_R \lesssim 10^7$ GeV obtained in the considered CMSSM scenarios for a standard cosmological history strongly disfavors thermal leptogenesis. However, if entropy is released after NLSP decoupling, a dilution factor of $\Delta \simeq 10^4$ can render thermal leptogenesis viable for $T_R \simeq 10^{13}$ GeV.
Standard thermal leptogenesis usually requires $T_R \gtrsim 10^9$ GeV \cite{17}. However, late-time entropy production dilutes the baryon asymmetry which is generated well before NLSP decoupling,

$$\eta(T_{\text{after}}) = \frac{1}{\Delta} \eta(T_{\text{before}}).$$

Therefore, the baryon asymmetry before entropy production must be larger by a factor of $\Delta$ in order to compensate for the dilution. For $\Delta \approx 10^3$, this can be achieved in the case of hierarchical neutrinos for $M_{R1} \sim T_R \approx 10^{13}$ GeV, as can be seen in Fig. 7 (a) of Ref. \cite{61} and in Fig. 2 of Ref. \cite{62}. Here $M_{R1}$ is the mass of the lightest among the heavy right-handed Majorana neutrinos.

In Fig. 6 the dotted (blue in the web version) lines show a scenario in which a dilution factor of $\Delta = 10^4$ is generated in the out-of-equilibrium decay of a heavy particle $X$. Because of $\rho_X(10 \text{ GeV}) = 8 \rho_{\text{had}}(10 \text{ GeV})$, the Hubble rate can be enhanced already during the decoupling phase of the NLSP, which leads to an increase of $T^{\text{NLSP}}_\text{NLSP}$ and $Y_{\text{NLSP}}(T^{\text{NLSP}}_\text{NLSP})$. In the results shown below, we account for this by using a modified version of the micrOMEGAs code.\(^{12}\) After entropy production, the net effect is still a significant reduction of $Y_{\text{NLSP}}(T_0)$. For the same initial conditions, $\Delta = 2 \times 10^4$—and thereby an additional reduction of $Y_{\text{NLSP}}(T_0)$ by a factor of two—can be achieved by lowering $T_{\text{after}}$ from 4.9 MeV down to 2.5 MeV.

We consider these two scenarios for $\tan \beta = 30$, $A_0 = 0$, $\mu > 0$, and $m_{\tilde{\nu}_L} = m_{\nu_L}$, in Fig. 8. Here the shaded (green in the web version) bands indicate the region in which $0.075 \leq \Omega m_0 h^2 \leq 0.126$ for $T_R = 10^{13}$ GeV and $\Delta = 10^4$ (dark) and $2 \times 10^6$ (medium). In addition, the corresponding evolution of the $^6\text{Li}$ bound is shown by the dot-dashed (red in the web version) lines. For $\Delta = 10^4$, the regions below the associated two rightmost curves and to the right of the associated leftmost curve are allowed. For $\Delta = 2 \times 10^4$, the cosmologically allowed region is the $^7\text{Li}$ NLSP region below the line labeled accordingly. The gray regions are identical to the ones in Fig. 5.

We find that the $^6\text{Li}$ bound cannot be evaded for the $\tan \beta = 10$ scenarios even for $\Delta = 2 \times 10^4$ since $Y_{\text{NLSP}}(T_0)$ becomes larger. However, the $^6\text{Li}$ bound given in Fig. 4 of Ref. \cite{22} depends linearly on the assumed limiting primordial abundance \cite{19} that is subject to uncertainties; cf. Ref. \cite{50}. Accordingly, for a limiting abundance that is a factor of two above the value given in \cite{19}, one obtains the $^6\text{Li}$ bound labeled with $\Delta = 2 \times 10^4$ in Fig. 8 for the scenario with $\tan \beta = 30$ and $\Delta = 10^4$.

We thank M. Pospelov for bringing this point to our attention.

\footnote{\label{Footnote1}The $Y_{\text{NLSP}}$ contours shown in Fig. 8 do not apply in this section.}

\footnote{\label{Footnote2}We thank M. Pospelov for bringing this point to our attention.}

CONCLUSION

Using the full gauge-invariant result for $\Omega_{\text{G}}^{\text{TP}}$ to leading order in the Standard Model gauge couplings \cite{11}, we have studied bounds on $T_R$ from the constraint $\Omega_{\tilde{\nu}} \leq \Omega_{\text{had}}$. Our results take into account the dependence of $\Omega_{\text{G}}^{\text{TP}}$ on the masses of the gauginos associated with the

\begin{figure}[h]
\centering
\includegraphics[width=\textwidth]{fig8.png}
\caption{The effect of entropy production after NLSP decoupling for $T_R = 10^{13}$ GeV and $\Delta \geq 10^4$ in the $(m_{1/2}, m_{\tilde{\nu}_L})$ plane for $\tan \beta = 30$, $A_0 = 0$, $\mu > 0$, and $m_{\tilde{\nu}_L} = m_{\nu_L}$. The shaded (green in the web version) bands show the region in which even slightly higher values of $T_R$ are possible for $\Delta = 10$. For $T_R = 10^{13}$ GeV and $\Delta \geq 10^4$, the generated baryon asymmetry is diluted too strongly in order to explain the observed baryon asymmetry. Such scenarios are particularly promising since the long-lived $^7\text{Li}$ NLSP could provide striking signatures of gravitino dark matter at future colliders.}
\end{figure}
Standard Model gauge group SU(3)c × SU(2)L × U(1)Y. This has allowed us to explore the dependence of the T_R bounds on the gaugino-mass relation at the scale of grand unification M_GUT.

Within the CMSSM, we have explored gravitino dark matter scenarios and the associated T_R bounds for m_{\tilde{G}} \lesssim 1 \text{ GeV} and for temperatures as low as 10^7 \text{ GeV}. Taking into account the restrictive constraint from bound-state effects of long-lived negatively charged staus on the primordial ^6\text{Li} abundance 22, we find that T_R \lesssim 10^7 \text{ GeV} is the highest cosmologically viable temperature of the radiation-dominated epoch in case of a standard thermal history of the Universe. This imposes a serious constraint on model building for inflation. Moreover, thermal leptogenesis seems to be strongly disfavored in the considered regions of the CMSSM parameter space.

With late-time entropy release, the obtained limit T_R \lesssim 10^7 \text{ GeV} can be relaxed. For example, the dilution of the thermally produced gravitino yield by a factor of 10 relaxes the T_R bound by about one order of magnitude in regions where \Omega_{\tilde{G}}^\text{T} dominates \Omega_{\tilde{G}}. In the case of entropy production after NLSP decoupling, the yield of the NLSP prior to its decay, Y_{\text{NLSP}} , is reduced so that the BBN constraints can be weakened. Although the ^6\text{Li} bound is persistent, we find that it disappears provided Y_{\text{NLSP}} is diluted by a factor of $\Delta \gtrsim 10^4$.

We have discussed the viability of thermal leptogenesis in a cosmological scenario with entropy production after NLSP decoupling. We find that successful thermal leptogenesis can be revived in generic regions of the CMSSM parameters space for $M_{\tilde{R}1} \sim T_R \simeq 10^{13} \text{ GeV}$ and $\Delta \gtrsim 10^4$, where $M_{\tilde{R}1}$ is the mass of the lightest among the heavy right-handed Majorana neutrinos.

Remarkably, for a dilution factor of $\Delta \gtrsim 10^4$, the $\tilde{\tau}_1$ NLSP region with $m_{\tilde{\tau}_1} \lesssim 200 \text{ GeV}$ reopens as a cosmologically allowed region in the CMSSM with the gravitino LSP. A long-lived $\tilde{\tau}_1$ in this mass range could provide striking signatures of gravitino dark matter at future colliders 63, 64, 65, 66.

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