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Gravitino Dark Matter from Inflaton Decay

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Abstract

We discuss a scenario that gravitinos produced non-thermally by an inflaton decay constitute dark matter in the present universe. We find that this scenario is realized for wide ranges of the inflaton mass and the vacuum expectation value. What is intriguing about this scenario is that the gravitino dark matter can have a relatively large free streaming length at matter-radiation equality, which can be probed by future observation on QSO-galaxy strong lens system.

In spite of accumulating observational data supporting the presence of the dark matter (DM) in our universe [1], we have not yet identified what DM is made of. Among many candidates proposed thus far, the gravitino, a supersymmetric parter of the graviton, is particularly interesting, and it has been thoroughly investigated in connection with leptogenesis [2] and the collider signatures [3].

The gravitinos are copiously produced by particle scattering in thermal plasma, once the decay of the inflaton reheats the universe. If the gravitino is the lightest supersymmetric particle (LSP), it is stable and can be a good candidate for DM [4, 5, 6, 7, 8]. The gravitino abundance is directly related to the reheating temperature, T_R. In particular, for the gravitino mass $m_{3/2}$ larger than $O(10)$GeV, the required reheating temperature for the gravitinos to be DM is so high, $T_R \gtrsim 10^9$GeV [9, 10], that the thermal leptogenesis scenario may work [11].

However, the detailed studies on the big bang nucleosynthesis (BBN) revealed that the abundance and the lifetime of the next-to-lightest supersymmetric particle (NLSP) are tightly constrained [12]. This drove the above attractive scenario into a corner, since the lifetime of NLSP tends to be longer than the BBN bound especially for $m_{3/2}$ larger than $O(10)$GeV, signaling the need for some changes. Several solutions has been proposed; e.g., a late-time entropy production [13, 14, 15] and a theory with R-parity violation [16, 17]. Another is to abandon thermal leptogenesis and consider a non-thermal leptogenesis scenario [18, 19] instead, which requires a lower reheating temperature, $T_R \gtrsim 10^6$GeV. Then the gravitino can account for the observed relic density even for $m_{3/2}$ lighter than $O(10)$GeV, making it easier to evade the constraints from the NLSP decay. One drawback of this approach however is that one needs to introduce ad hoc couplings of the inflaton with the right-handed neutrinos.

Furthermore, it has been recently pointed out that the gravitinos are generically produced by an inflaton decay [20, 21, 22, 23, 24, 25, 26]. Since such non-thermal gravitino production generically occurs for the most inflation models, it is worth studying how it affects the conventional picture on the gravitino DM scenario. In this letter, we pursue a possibility that the gravitinos produced by an inflaton decay constitute a dominant component of DM. Generically, the required reheating temperature becomes lower than without the non-thermal production. This makes it difficult to integrate the thermal leptogenesis scenario into this framework. As we will see later, however, the right-handed (s)neutrinos are generated by the inflaton decay, and their subsequent decay may generate a right amount of the baryon asymmetry via leptogenesis for certain values of the inflaton parameters [27]. What is particularly appealing about this scenario is that both the gravitino DM and the non-thermal leptogenesis can be realized without introducing any couplings ad hoc by hand, if the inflaton parameters satisfy certain conditions. In addition, the produced gravitinos can have a large velocity at matter-radiation equality, which affects the growth of density fluctuations of DM [a]. Future observations on e.g. QSO-galaxy strong lens system [32] may be able to support or refute this scenario.

Let us first briefly review the recent development on the gravitino production from the inflaton decay. There are three gravitino production processes; (a) the gravitino pair production [20, 21, 22, 23]; (b) spontaneous decay at tree level [24]; (c) anomaly-induced decay at one-loop level [25]. For the processes listed above, the gravitino production rate can be expressed as

$$\Gamma_{3/2} = \frac{x}{32\pi} \left(\frac{\langle \phi \rangle}{M_P} \right)^2 \frac{m_\phi^3}{M_P^2}, \tag{1}$$

where m_ϕ is the inflaton mass, $\langle \phi \rangle$ a vacuum expectation value (VEV) of the inflaton, and $M_P = 2.4 \times 10^{18}$GeV the reduced Planck mass. Here it should be noted that $\langle \phi \rangle$ is evaluated at the potential minimum after inflation. The precise value of the numerical coefficient x depends on the production processes, possible non-renormalizable couplings in the Kähler potential, and the detailed structure of the supersymmetry (SUSY) breaking sector [26]. To be concrete, let us assume the minimal Kähler potential and the dynamical SUSY breaking (DSB) [33] with a dynamical scale Λ. In the DSB scenario, the SUSY breaking field z can acquire a large mass m_z, which is assumed to be roughly equal to the dynamical scale $\Lambda \sim \sqrt{m_{3/2} M_P}$ in the following. Such a simplification does not essentially change our arguments. For a low-inflation model with $m_\phi < \Lambda$, the process (a) becomes effective, and $x = 1$. On the other hand, for the inflaton mass larger than Λ, the processes (b) and (c) become effective instead. The inflaton decays into the hidden quarks in the SUSY breaking sector via Yukawa couplings (process (b)), or into the hidden gauge sector via anomalies (process (c)). Since the hidden quarks and gauge bosons (and gauginos) are energetic when they are produced, they are expected to form jets and produce hidden hadrons through the

[a] Such a DM candidate with a large velocity and its astrophysical implication was first discussed in Refs .[28, 29, 30], and intensively studied in connection with the so-called superWIMP mechanism [31].

strong gauge interactions. The gravitinos are likely generated by the decays of the hidden hadrons as well as in the cascade decay processes in jets. We denote the averaged number of the gravitinos produced per each jet as $N_{3/2}$. Then x is given by [26] [b]

$$x \;=\; \frac{N_{3/2}}{8\pi^2}\left(\frac{1}{2}N_y|Y_h^2| + N_g\alpha_h^2(T_g^{(h)} - T_r^{(h)})^2\right), \tag{2}$$

where Y_h and α_h are the Yukawa coupling and a fine structure constant of the hidden gauge group, respectively, N_y denotes a number of the final states for the process (b), N_g is a number of the generators of the gauge group, and $T_g^{(h)}$ and $T_r^{(h)}$ are the Dynkin indices of the adjoint representation and the matter fields in the representation r. Although x depends on the structure of the SUSY breaking sector, its typical magnitude is $O(10^{-3} - 10^{-2})$ for $m_\phi > \Lambda$ [c]. To be concrete we take

$$x = \left\{ \begin{array}{ccc} 1 & \text{for} & m_\phi < \Lambda \\ 10^{-3} \;\text{ or }\; 10^{-2} & \text{for} & m_\phi > \Lambda \end{array} \right., \tag{3}$$

in the following.

Using the gravitino production rate given above, we can estimate the abundance of the gravitinos non-thermally produced by an inflaton decay:

$$\begin{aligned} Y_{3/2}^{(NT)} &= 2\frac{\Gamma_{3/2}}{\Gamma_\phi}\frac{3T_R}{4m_\phi}, \\ &\simeq 7\times10^{-11}\,x\left(\frac{g_*}{200}\right)^{-\frac{1}{2}}\left(\frac{\langle\phi\rangle}{10^{15}\mathrm{GeV}}\right)^2\left(\frac{m_\phi}{10^{12}\mathrm{GeV}}\right)^2\left(\frac{T_R}{10^6\mathrm{GeV}}\right)^{-1}, \end{aligned} \tag{4}$$

where g_* counts the relativistic degrees of freedom, and Γ_ϕ denotes the total decay rate of the inflaton that is related to the reheating temperature as

$$\Gamma_\phi \;\equiv\; \left(\frac{\pi^2 g_*}{10}\right)^{\frac{1}{2}}\frac{T_R^2}{M_P}. \tag{5}$$

Equivalently, the gravitino density parameter is

$$\Omega_{3/2}^{(NT)}h^2 \;\simeq\; 0.02\,x\left(\frac{g_*}{200}\right)^{-\frac{1}{2}}\left(\frac{m_{3/2}}{1\mathrm{GeV}}\right)\left(\frac{\langle\phi\rangle}{10^{15}\mathrm{GeV}}\right)^2\left(\frac{m_\phi}{10^{12}\mathrm{GeV}}\right)^2\left(\frac{T_R}{10^6\mathrm{GeV}}\right)^{-1}, \tag{6}$$

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[b] If the Kähler potential takes a form of the sequestered type, the spontaneous decay through Yukawa couplings is suppressed [24, 25].

[c] Roughly, we expect $N_{3/2} = O(1 - 10^2)$, $N_g = O(1)$, $\alpha_h = O(0.1)$, and $T_g^{(h)} - T_r^{(h)} = O(1)$, while Y_h strongly depends on the SUSY breaking models. Note also that the gravitino can be produced through the Yukawa interaction in the messenger sector, if the inflaton mass is larger than the messenger scale.

where h is the present Hubble parameter in units of 100km/s/Mpc. Note that the gravitino abundance is inversely proportional to the reheating temperature. Due to this feature, the non-thermal gravitino production tends to require a relatively low T_R to realize the gravitino DM scenario. Indeed, by solving $\Omega_{3/2}^{(NT)} h^2 = 0.11$ [36] with respect to T_R, we obtain

$$T_R \;\simeq\; 2 \times 10^5 \,\text{GeV}\, x \left(\frac{g_*}{200}\right)^{-\frac{1}{2}} \left(\frac{m_{3/2}}{1\text{GeV}}\right) \left(\frac{\langle\phi\rangle}{10^{15}\text{GeV}}\right)^2 \left(\frac{m_\phi}{10^{12}\text{GeV}}\right)^2 . \tag{7}$$

So, if T_R is given by the above value, the non-thermally produced gravitino has a right abundance to become a dominant component of DM. Generically, one has to introduce a coupling of the inflaton to the standard-model sector with an appropriate strength, in order to realize T_R given by Eq. (7). However, there is a natural way to induce the reheating, and we will discuss this possibility later.

For the non-thermally produced gravitinos to account for DM, several conditions must be met. First, the gravitino production by thermal scatterings should give only negligible contribution to the DM abundance. The abundance of the gravitinos produced by thermal scatterings is given by [6, 34, 35]

$$\Omega_{3/2}^{(th)} h^2 \;\simeq\; 0.14 \left(\frac{m_{\tilde{g}_3}}{300\text{GeV}}\right)^2 \left(\frac{m_{3/2}}{1\text{GeV}}\right)^{-1} \left(\frac{T_\text{R}}{10^8\ \text{GeV}}\right), \tag{8}$$

where $m_{\tilde{g}_3}$ is the gluino running mass evaluated at $T = T_R$. Requiring $\Omega_{3/2}^{(th)} h^2$ to be less than the observed DM abundance, $\Omega_{DM} h^2 \simeq 0.11$, T_R is bounded above:

$$T_R \;\lesssim\; 8 \times 10^7 \,\text{GeV} \left(\frac{m_{\tilde{g}_3}}{300\text{GeV}}\right)^{-2} \left(\frac{m_{3/2}}{1\text{GeV}}\right) . \tag{9}$$

This constraint is valid for $m_{3/2} \gtrsim 100\,\text{keV}$, which is satisfied for the parameter space concerned as shown later.

Another constraint comes from the recent discovery that, once the inflaton acquires a non-vanishing VEV, the inflaton decays into the visible sector through the top Yukawa coupling [24]. Due to the presence of this decay process, T_R cannot be arbitrarily low. Indeed, it is bounded below as

$$T_R \;\gtrsim\; 1.9 \times 10^3 \,\text{GeV}\, |Y_t| \left(\frac{g_*}{200}\right)^{-\frac{1}{4}} \left(\frac{\langle\phi\rangle}{10^{15}\,\text{GeV}}\right) \left(\frac{m_\phi}{10^{12}\,\text{GeV}}\right)^{\frac{3}{2}} , \tag{10}$$

where Y_t is the top Yukawa coupling. The inequality is saturated if the inflaton has no direct couplings with any other fields in superpotential [d].

[d] Note that we assume the minimal Kähler potential in the Einstein frame.

The last constraint arises from the fact that the non-thermally produced gravitinos can have a large velocity at matter-radiation equality, in contrast to the gravitinos produced by thermal scatterings. This not only limits the parameter space, but also provides a possibility that the scenario may be probed by future observation on QSO-galaxy strong lens system. Let us estimate the comoving free streaming length of the gravitino at matter-radiation equality, assuming that it has an initial energy, $\epsilon m_\phi/2$ when produced. For $m_\phi < \Lambda$, we have $\epsilon = 1$ since a pair of the gravitinos is directly produced by the inflaton decay. On the other hand, for $m_\phi > \Lambda$, multiple gravitinos are indirectly generated by the inflaton decay, and so, its energy tends to be smaller than $m_\phi/2$; we expect $\epsilon \lesssim O(N_{3/2}^{-1}) = O(10^{-3} - 0.1)$. To be concrete we will take $\epsilon = 10^{-3}$ or 10^{-2} for $m_\phi > \Lambda$. The comoving free streaming length λ_{FS} at matter-radiation equality is defined by

$$\lambda_{FS} \equiv \int_{t_D}^{t_{\rm eq}} \frac{v_{3/2}(t)}{a(t)} dt, \tag{11}$$

where $a(t)$ is the scale factor, and t_D and $t_{\rm eq}(\sim 2\times 10^{12}\,{\rm sec})$ denote the time at the inflaton decay and at matter-radiation equality, respectively. $v_{3/2}$ is the velocity of the gravitino, given by

$$v_{3/2}(t) = \frac{|\mathbf{p}_{3/2}|}{E_{3/2}} \simeq \frac{\frac{\epsilon m_\phi}{2}\left(\frac{a_D}{a(t)}\right)}{\sqrt{m_{3/2}^2 + \frac{\epsilon^2 m_\phi^2}{4}\left(\frac{a_D}{a(t)}\right)^2}}, \tag{12}$$

where we have approximated $m_\phi \gg m_{3/2}$, and a_D is the scale factor at the inflaton decay. Integrating (11) yields

$$\begin{aligned} \lambda_{FS} &\simeq \frac{1}{H_0\sqrt{1+z_{\rm eq}}} X^{-1} \sinh^{-1} X, \\ &\sim 0.09\,{\rm Mpc}\,\epsilon \ln(2X) \left(\frac{g_*}{200}\right)^{-\frac{1}{4}} \left(\frac{m_{3/2}}{1{\rm GeV}}\right)^{-1} \left(\frac{m_\phi}{10^{12}{\rm GeV}}\right) \left(\frac{T_R}{10^5{\rm GeV}}\right)^{-1} \end{aligned} \tag{13}$$

with

$$\begin{aligned} X &\equiv \frac{2m_{3/2}}{\epsilon m_\phi}\frac{a_{\rm eq}}{a_D}, \\ &\simeq 8\times 10^2\,\epsilon^{-1} \left(\frac{g_*}{200}\right)^{-\frac{1}{4}} \left(\frac{m_{3/2}}{1{\rm GeV}}\right) \left(\frac{m_\phi}{10^{12}{\rm GeV}}\right)^{-1} \left(\frac{T_R}{10^5{\rm GeV}}\right), \end{aligned} \tag{14}$$

where H_0 is the Hubble parameter at present, and $z_{\rm eq}$ and $a_{\rm eq}$ are the red-shift and the scale factor at the matter-radiation equality. In the second equation of (13), we have assumed $X \gg 1$ and used $H_0^{-1} \sim 4\times 10^3\,{\rm Mpc}$ and $z_{\rm eq} \sim 3000$. In Eq. (14), we have used $a_D/a_{\rm eq} =$

$(\Gamma_\phi \cdot t_{\rm eq})^{-1/2}$. The constraint from Ly-α clouds, $\lambda_{FS} \lesssim 1\,{\rm Mpc}$, implies $X \gtrsim 450$. We therefore obtain a constraint on T_R as

$$T_R \;\gtrsim\; 5\times 10^4\,{\rm GeV}\,\epsilon \left(\frac{g_*}{200}\right)^{-\frac{1}{4}} \left(\frac{m_{3/2}}{1{\rm GeV}}\right)^{-1} \left(\frac{m_\phi}{10^{12}{\rm GeV}}\right). \tag{15}$$

The meaning of this constraint is clear: the reheating must occur so early that the velocity of the produced gravitino becomes small enough due to redshift by the matter-radiation equality.

Thus, if the reheating temperature T_R is given by (7) and satisfies the above constraints (9), (10), and (15) in addition to the BBN constraint $T_R \gtrsim 10\,{\rm MeV}$ [37, 38, 39], the non-thermally produced gravitinos account for DM. As mentioned above, one may have to add appropriate couplings of the inflaton to light degrees of freedom, in order to realize T_R given by (7). However there is one interesting possibility that the reheating is induced by the decay through the top Yukawa coupling. Then the inequality (10) becomes saturated. This is the case if there are no direct couplings of the inflaton with any other fields in the superpotential. The presence of the decay process through the top Yukawa coupling not only constrains the reheating temperature, but also provides an intriguing way to induce the reheating. For the moment let us pursue this possibility. From (7) and (10), we obtain

$$\left(\frac{\langle\phi\rangle}{10^{15}{\rm GeV}}\right)\left(\frac{m_\phi}{10^{12}{\rm GeV}}\right)^{\frac{1}{2}} \;\simeq\; 0.01\,|Y_t| x^{-1} \left(\frac{g_*}{200}\right)^{\frac{1}{4}} \left(\frac{m_{3/2}}{1{\rm GeV}}\right)^{-1}. \tag{16}$$

Thus, if the inflaton parameters, m_ϕ and $\langle\phi\rangle$, satisfy the above relation (16), the non-thermally produced gravitino has a just right abundance to be DM. Interestingly, the free streaming length becomes independent of the inflaton parameters and the gravitino mass in this case. Indeed, λ_{FS} is approximately given by

$$\lambda_{FS} \;\simeq\; 1\times 10^2\,{\rm kpc} \left(\frac{g_*}{200}\right)^{-\frac{1}{4}} \left(\frac{|Y_t|}{0.6}\right)^{-2} \left(\frac{\epsilon\, x}{10^{-5}}\right), \tag{17}$$

The Ly-α constraint requires $\epsilon\, x \lesssim 10^{-4}$, which is naturally satisfied for a high-scale inflation model with $m_\phi > \Lambda$. It is intriguing that the gravitino DM scenario points to a high-scale inflation model with $m_\phi > \Lambda$ and predict a relatively large free streaming length, as long as the reheating is induced by the top Yukawa coupling. For $x = 10^{-3} \sim 10^{-2}$ and $\epsilon = 10^{-3} \sim 10^{-1}$, the comoving free streaming length takes a value from 10 kpc up to 1 Mpc (limited by the Ly-α constraint).

Now let us consider the inflaton decay into the right-handed (s)neutrinos thorough large Majorana mass terms:

$$W = \frac{M_i}{2} N_i N_i, \tag{18}$$

where $i = 1, 2, 3$ is the family index. We consider the inflaton decay into the lightest right-handed (s)neutrino N_1 for simplicity, assuming that the decay into the heavier ones, N_2 and N_3, are kinematically forbidden. We drop the family index in the following. The partial decay rate of the inflaton into the right-handed (s)neutrinos is [cf. [24]]

$$\Gamma_N \simeq \frac{1}{16\pi} \left(\frac{\langle \phi \rangle}{M_P} \right)^2 \frac{m_\phi M^2}{M_P^2} \sqrt{1 - \frac{4M^2}{m_\phi^2}}, \tag{19}$$

where we have taken account of both the decay into the right-handed neutrinos and that into the right-handed sneutrinos. Note that one does not have to introduce any direct couplings of the inflaton with the right–handed neutrinos to induce the decay. The decay proceeds as long as the inflaton acquires a nonzero VEV.

The lepton asymmetry can be produced by the decay of the right-handed (s)neutrinos, if CP is violated in the neutrino Yukawa matrix [2]. The resultant lepton asymmetry is given by

$$\frac{n_L}{s} \simeq \frac{3}{2} \epsilon_1 \, B_N \frac{T_R}{m_\phi}, \tag{20}$$

where $B_N \equiv \Gamma_N / \Gamma_\phi$ denotes the branching ratio of the inflaton decay into the (s)neutrinos. The asymmetry parameter ϵ_1 is given by [2, 40]

$$\epsilon_1 \simeq 2.0 \times 10^{-10} \left(\frac{M}{10^6 \mathrm{GeV}} \right) \left(\frac{m_{\nu_3}}{0.05 \mathrm{eV}} \right) \delta_{\rm eff}, \tag{21}$$

where m_{ν_3} is the heaviest neutrino mass and $\delta_{\rm eff} \leq 1$ represents the effective CP-violating phase. The baryon asymmetry is obtained via the sphaleron effect: [41]

$$\frac{n_B}{s} = -\frac{8}{23} \frac{n_L}{s}. \tag{22}$$

Using the above relations, we obtain the right amount of baryon asymmetry,

$$\begin{aligned} \frac{n_B}{s} &\simeq 1 \times 10^{-9} \left(\frac{g_*}{200} \right)^{-\frac{1}{2}} \left(\frac{M}{10^{13} \mathrm{GeV}} \right)^3 \left(\frac{\langle \phi \rangle}{10^{16} \mathrm{GeV}} \right)^2 \left(\frac{T_R}{10^6 \mathrm{GeV}} \right)^{-1} \left(\frac{m_{\nu_3}}{0.05 \mathrm{eV}} \right) \delta_{\rm eff}, \\ &\simeq 5 \times 10^{-11} \left(\frac{g_*}{200} \right)^{-\frac{1}{4}} \left(\frac{M}{10^{13} \mathrm{GeV}} \right)^3 \left(\frac{\langle \phi \rangle}{10^{16} \mathrm{GeV}} \right) \left(\frac{m_\phi}{10^{14} \mathrm{GeV}} \right)^{-\frac{3}{2}} \left(\frac{m_{\nu_3}}{0.05 \mathrm{eV}} \right) \delta_{\rm eff}, \end{aligned} \tag{23}$$

where we have assumed that the inequality (10) is saturated in the second equality. Note that M cannot exceed $m_\phi/2$. Therefore, the baryon asymmetry is proportional to positive powers of m_ϕ, if M is set to be a value that maximizes the asymmetry.

In Figs. 1 and 2, we show the parameter space where the reheating temperature (7) satisfies the above constraints (9), (10), and (15), in addition to the BBN constraint $T_R \gtrsim 10\,\mathrm{MeV}$. In the shaded (blue) regions, the baryon asymmetry can be explained by the non-thermal leptogenesis scenario discussed above, if an appropriate value of $M(\lesssim 10^{15}\mathrm{GeV})$ is chosen (we set $m_{\nu_3} = 0.05\,\mathrm{eV}$ and $\delta_{\rm eff} = 1$). We have chosen several values of the gravitino mass: $m_{3/2} = 100\,\mathrm{MeV}$, $1\,\mathrm{GeV}$, and $10\,\mathrm{GeV}$ [e]. For smaller $m_{3/2}$, one needs to generate more gravitinos, due to which the allowed region shifts upward. At the same time, the constraints from (9) and (15) become severer for smaller $m_{3/2}$, reducing the allowed space. Thus, if $x = 10^{-3}(10^{-2})$ for $m_\phi > \Lambda$, the gravitino mass should be larger than 1 MeV (100 keV) for the non-thermally produced gravitinos to account for DM, since otherwise there is no allowed region for $\langle\phi\rangle \lesssim M_P$. If x becomes larger for $m_\phi > \Lambda$, the allowed region shifts downward, and a smaller value of the gravitinos mass becomes allowed. On the other hand, if x becomes smaller due to e.g. conformal sequestering [26], we have more parameter space for the non-thermal leptogenesis to work successfully.

The dotted (red) lines correspond to the special case that the reheating is solely induced by the decay through the top Yukawa coupling and the non-thermally produced gravitinos become DM. Therefore the inflaton parameters on the dotted (red) lines are particularly interesting in a sense that one does not have to introduce any couplings ad hoc by hand; the decay spontaneous proceeds through the top Yukawa coupling, and the gravitino has just a right abundance to become DM. Note that the free streaming length is constant along the dotted (red) lines and independent of $m_{3/2}$, m_ϕ and $\langle\phi\rangle$, as mentioned before.

In Figs. 1 and 2, we also show the contours of the free streaming length $\lambda_{FS} = 1\,\mathrm{kpc}$, 10 kpc, 100 kpc, and 1 Mpc. The future submillilensing observations can cover $\lambda_{FS} \gtrsim 2$ kpc [32]. In particular, since the interesting case that the reheating occurs through the top Yukawa coupling (dotted red lines) predicts the gravitino DM with a relatively large free streaming length ($\gtrsim 10$ kpc), it can be probed by future observations. Such a large free

[e] Note that we have not taken account of the constraints on the NSLP decay. Therefore, the figure in the case of $m_{3/2} = 10\,\mathrm{GeV}$ is valid only if the cosmological problems associated with the NLSP are somehow avoided by e.g. introducing R-parity violating operators with an appropriate magnitude [16, 17].

streaming length may also solve the missing satellite problem [45] and the cusp problem [46].

From the figures, one can see that relatively broad ranges of the inflaton mass and VEV are allowed. In particular, when combined with the non-thermal leptogenesis scenario, we are led to a high-scale inflation model with $m_\phi > \Lambda$. However, studying the parameter spaces of the representative high-scale inflation models (such as the hybrid [42] and smooth hybrid [43], and chaotic [44] [f] inflation models) in detail, one finds that only small part of the parameter space actually overlaps with the region where the non-thermal leptogenesis works, especially if x takes a value on the high side $\sim 10^{-2}$; those inflation models tend to predict lighter m_ϕ and larger $\langle\phi\rangle$ compared to those favored by the non-thermal leptogenesis. See Fig. 3. Such a tension may be ameliorated if one assumes some mechanism (e.g. the conformal sequestering) to suppress x to a smaller value.

In summary, we have considered a scenario that the non-thermally produced gravitinos from the inflaton decay become a dominant component of DM. Interestingly, if the reheating is induced solely by the decay through the top Yukawa coupling, a high-scale inflation model is required for the non-thermally produced gravitinos to account for DM, and the free streaming length λ_{FS} is predicted to be in the range between $O(10)$ kpc and $O(0.1)$ Mpc, independently of the inflaton parameters and the gravitino mass. Such large free streaming length may affect the growth of the density fluctuations in DM. The suppression of the density contrast below the free streaming scale results in the absence of the sub-halos. This feature may be supported or refuted by future observations on the QSO-galaxy strong lens system [32].

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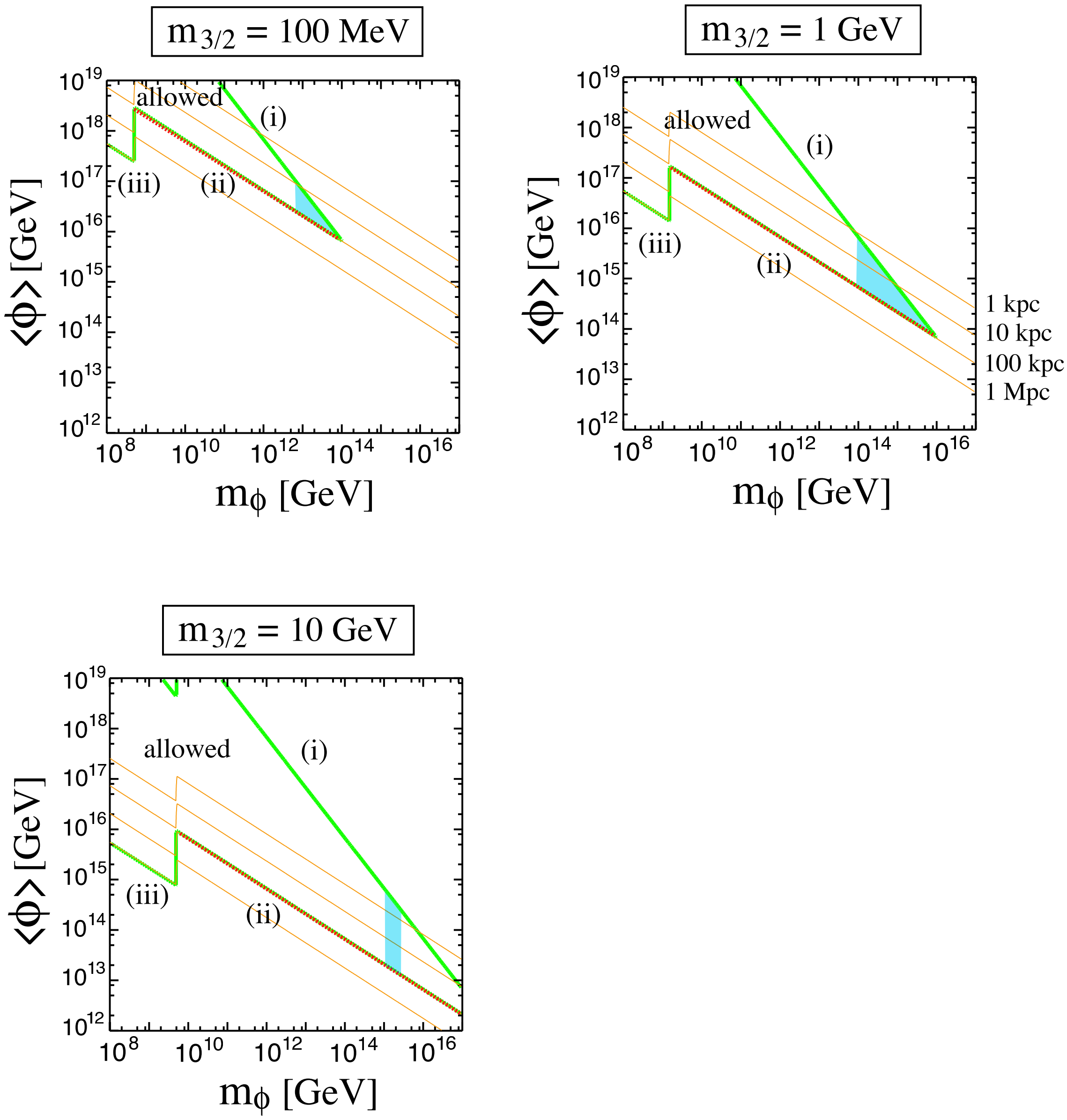

FIG. 1: In the regions surrounded by the solid (green) lines, the non-thermally produced gravitinos can account for the observed DM density, if T_R is given by (7). We have imposed the constraints from (i) thermal production of the gravitino (see (9)); (ii) decay through the top Yukawa coupling (see (10)); (iii) Ly-α clouds (see (15)). On the dotted (red) line, the reheating is induced solely by the decay via the top Yukawa coupling and the non-thermally produced gravitino explains DM. The thin solid (orange) lines are the contours of the free streaming length $\lambda_{FS} = 1\,$kpc, 10 kpc, 100 kpc, and 1 Mpc, from top to bottom. In the shaded (blue) regions the present baryon asymmetry can be explained by the non-thermal leptogenesis. We set $x = 10^{-3}$ and $\epsilon = 10^{-2}$ for $m_\phi > \Lambda$, respectively.

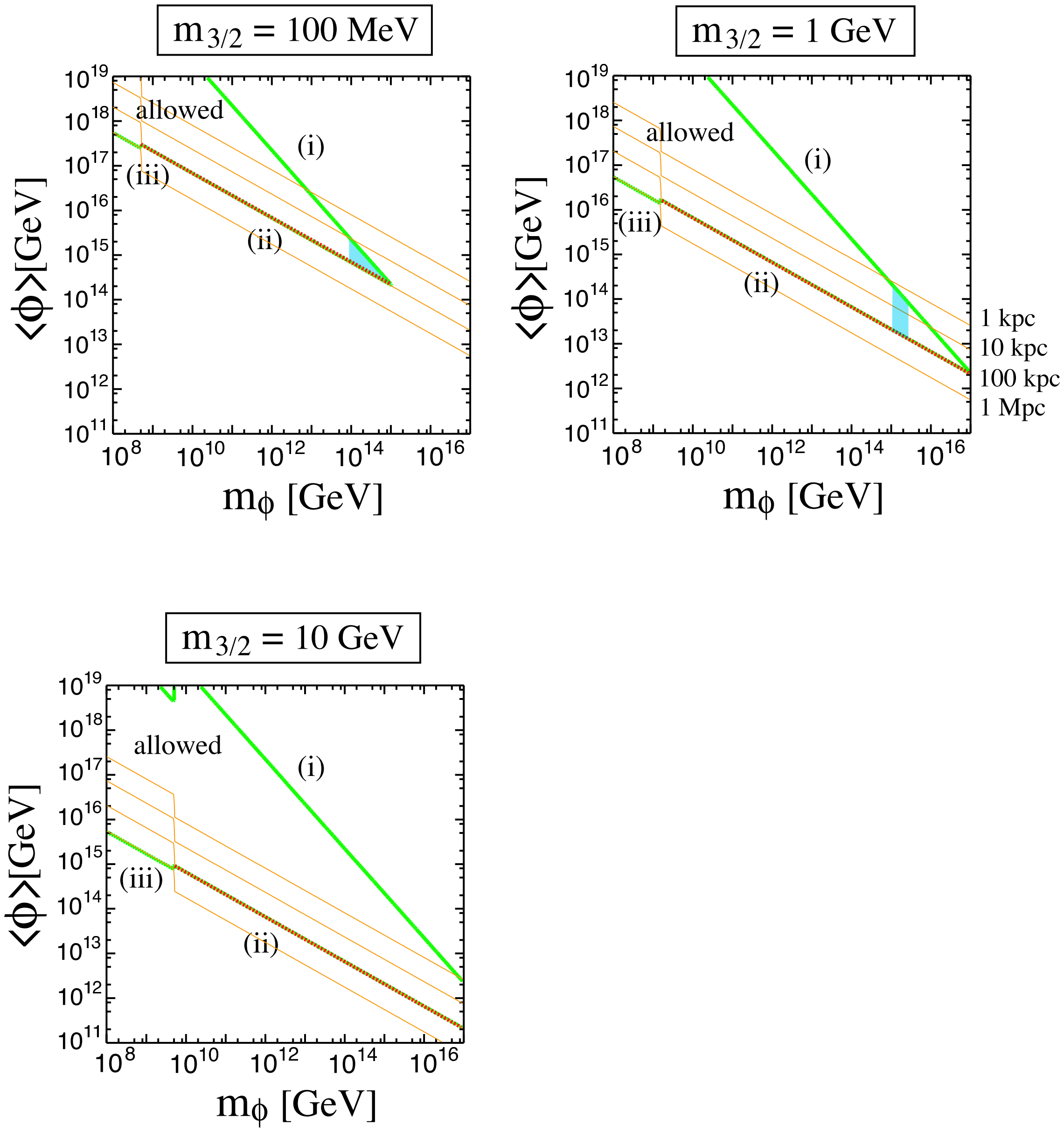

FIG. 2: Same as Fig. 1 except for $x = 10^{-2}$ and $\epsilon = 10^{-3}$ for $m_\phi > \Lambda$.

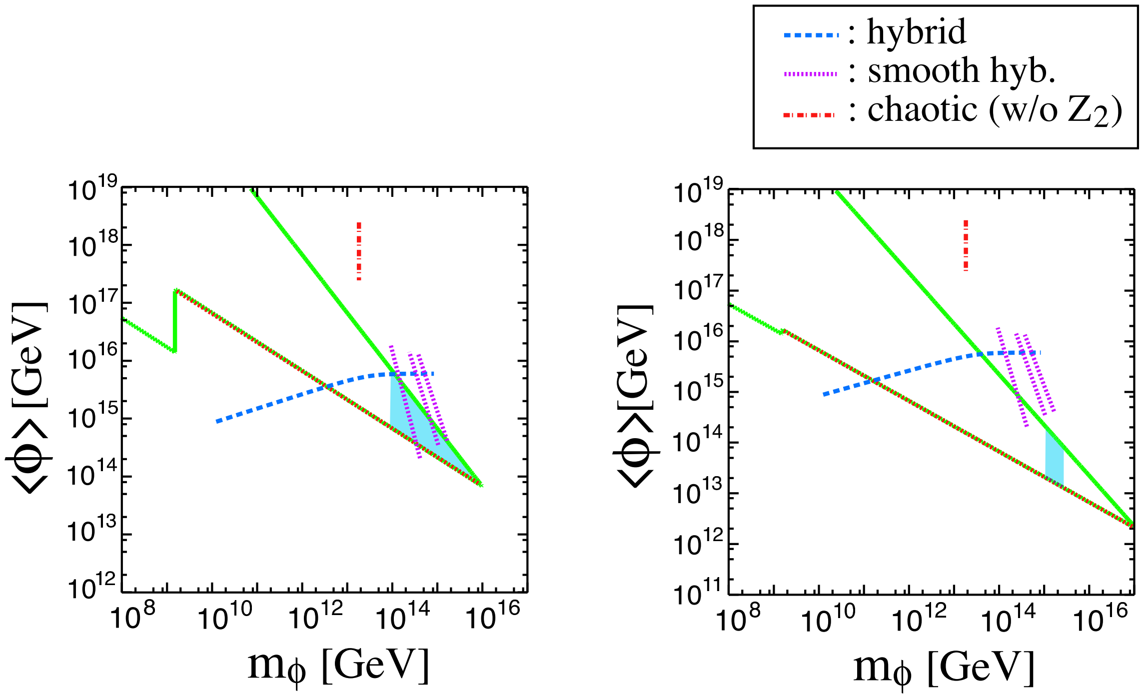

FIG. 3: We show the representative high-scale inflation models; hybrid [42] (thick long dashed (blue) line), smooth hybrid [43] (thick short dashed (purple) line), and chaotic [44] (long dashed dotted (red)) inflation models, superposed on the panels of $m_{3/2} = 1$GeV shown in Fig. 1 (left) and Fig. 2 (right).