A New Inflation Model with Anomaly-mediated Supersymmetry Breaking

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Abstract

If there are a large number of vacua, multi-inflation may be a more mediocre phenomenon rather than a single inflation. In the multi-inflation scenario, new inflation is most likely the last inflation, since its energy scale is naturally low. Furthermore, it may explain the observed spectral index of the cosmic microwave background radiations. We show, in this letter, that a new inflation model proposed in supergravity accounts for all the present observations assuming anomaly mediation of supersymmetry breaking. As a result, we find that the relic density of the winos is consistent with the observed dark matter density in a wide range of the wino mass, $100 \text{GeV} \lesssim m_{\tilde{w}} \lesssim 2 \text{ TeV}$, albeit for a low reheating temperature $T_R \simeq 10^{6-7} \text{ GeV}$. 
1 Introduction

The existence of a large number of vacua is the most exciting discovery of string theory \[1\]. This multiplicity of vacua called as "landscape of vacua" \[2\] provides a theoretical basis for the anthropic explanation on the small cosmological constant \[3\]. That is, if the cosmological constant, \(\Lambda_{\text{cos}}\), takes a wide range of values in the full vacua, one may find intelligent observers in sub-vacua where the \(\Lambda_{\text{cos}}\) takes a sufficiently small value for the presence of the observers. In this landscape the flat potential of a scaler field for inflation is also naturally explained, since it is necessary for the observers to exist. Furthermore, multi-inflation seems a more mediocre phenomenon in the landscape, rather than a single inflation \[4\] \[5\].

If the multi-inflation takes place, a lower-scale inflation starts at a later time. Thus, the last inflation we see today is most likely a low-scale inflation. We consider a new inflation as the last one, since new inflation is known to have naturally a low-scale Hubble constant. In a recent article \[6\] we have shown that a new inflation model \[7\] \[8\] constructed in the supergravity (SUGRA) predicts the spectral index \(n_s\) of the cosmic microwave background radiations as \(n_s \approx 0.95\) \[8, 9\] in a large parameter space. This result turns out to be well consistent with the recent WMAP observation \[10\], \(n_s = 0.951^{+0.015}_{-0.019}\) (68\%C.L.). Encouraged by this success, we examine, in the present letter, if this new inflation model is consistent with all other observations. We assume anomaly-mediation models for supersymmetry (SUSY) breaking, since gravity-mediation models suffer from a serious gravitino-overproduction problem \[11\] \[12\] \[13\] \[14\]. We stress, in particular, that the new inflation model (with anomaly-mediated SUSY breaking) may explain the observed dark matter density as well as the baryon asymmetry in the universe. As we show, the relic density of the wino LSP is consistent with the observed dark matter density in a wide range of the wino mass, \(100 \text{GeV} \lesssim m_{\tilde{w}} \lesssim 2 \text{TeV}\), albeit for a low reheating temperature \(T_R \approx 10^6-7\) GeV.

We also briefly discuss gauge-mediation models in the last section.
2 A new inflation model

Let us discuss a new inflation model considered in Ref. [7, 8]. In the model, the superpotential and the Kähler potential of an inflaton chiral superfield $\phi$ are given by

$$W_{\text{inf}} = v^2 \phi - \frac{g}{n+1} \frac{\phi^{n+1}}{M_G^{n-2}}, \quad (1)$$

and

$$K_{\text{inf}} = |\phi|^2 + \frac{k |\phi|^4}{4 M_G^2} + \cdots, \quad (2)$$

respectively. Here, $v^2$ denotes a dimensionful parameter and $g$ and $k$ denote dimensionless coupling constants. We take the parameters $v^2$ and $g$ positive without a loss of generality. And $n$ is an integer number greater than 2. Hereafter, we take the unit with the reduced Planck scale, $M_G \simeq 2.4 \times 10^{18}$ GeV, equal to one. The above superpotential is generic under a discrete $Z_{2n}$ $R$-symmetry with $\phi$’s charge 2. Then, the effective scalar potential of the inflaton $\varphi = \sqrt{2} \text{Re}[\phi]$ is well approximated by

$$V(\varphi) \simeq v^4 - \frac{k}{2} v^4 \varphi^2 - \frac{g}{2^{n-1}} v^2 \varphi^n + \frac{g^2}{2^n} \varphi^{2n}, \quad (3)$$

for the inflationary period near the origin, $\varphi = 0$. This potential is very flat for $n \geq 3$ and $|k| \ll 1$. In the following, we consider that the previous inflation drives the inflaton $\varphi$ to the origin [15], and assume $k > 0$ so that the inflaton $\varphi$ rolls down slowly to the potential minimum from near the origin. Note that the inflaton obtains a mass,

$$m_\varphi \simeq n g \phi_0^{n-1} \simeq n v^2 \left( \frac{v^2}{g} \right)^{-\frac{1}{n}}, \quad (4)$$

at the potential minimum,

$$\phi_0 = \frac{1}{\sqrt{2}} \varphi_0 \simeq \left( \frac{v^2}{g} \right)^{\frac{1}{n}}. \quad (5)$$

We should stress here that the present inflation model contains only four parameters, $v$, $k$, $g$ and $n$. We now show all parameters are determined by the observations, provided that the SUSY is broken at low energies, that is, the gravitino mass $m_{3/2} < 10^6$ GeV.

First of all, we should note one of the most remarkable features of the present new inflation model; the inflation scale $v$ is directly related to the gravitino mass [7]. The
important point is that a constant term in the superpotential is generated in the true minimum of the inflaton potential. Thus, the negative energy at the inflaton potential minimum is to be canceled out by the positive energy $\Lambda_{\text{SUSY}}^4$ induced by the SUSY breaking to have the cosmological constant vanishing. Thus, we have a condition as

$$\Lambda_{\text{SUSY}}^4 - 3|W_{\inf}(\phi_0)|^2 = 0. \quad (6)$$

Then, the gravitino mass is given by

$$m_{3/2} = W_{\inf}(\phi_0) \simeq \frac{nv^2}{n + 1} \left( \frac{v^2}{g} \right)^{\frac{1}{2}}. \quad (7)$$

We see that the gravitino mass is basically given by the inflation scale $v$.

Before determining the inflation scale $v$, we derive constraints on the parameters $n$ and $k$. The spectral index of the density fluctuations is given by [9, 6, 8]

$$n_s \simeq 1 - 6\epsilon + 2\eta, \quad (8)$$

$$\simeq 1 - 2k \left[ 1 + \frac{n-1}{1 + \frac{k}{n} (n-1)} e^{N_e k (n-2) - 1} \right], \quad (9)$$

where $\epsilon$ and $\eta$ denote so-called slow roll parameters at the horizon crossing, and $N_e$ is the e-folding number of the present universe [16]. Note that the spectral index depends on neither $v^2$ nor $g$ explicitly and is mainly given by the parameter $k$. As stressed in the introduction, the present new inflation model predicts the spectral index $n_s \simeq 0.95$ for $n \geq 4$ and $k \lesssim 10^{-2}$ (see Fig. [11]), which is well consistent with the recent WMAP result, $n_s = 0.951^{+0.015}_{-0.019}$ (68\% C.L.) [10]. The WMAP observation favors models with $n \geq 4$ and $k \lesssim 10^{-2}$. We, therefore, take $n \geq 4$ and $k \lesssim 10^{-2}$ in the following discussion.

We now come to the point to determine the inflation scale $v$. The inflation scale $v$ is given by

$$v \simeq \left( \frac{5}{6} \pi n \delta \right)^{\frac{n-2}{2n-6}} \left( \frac{1}{n(n-2)N_e} \right)^{\frac{2}{2n-6}} \left( \frac{1}{g} \right)^{\frac{1}{n-6}}, \quad (10)$$

for $n \geq 4$ and $k \lesssim 10^{-2}$, where we have neglected a weak dependence on $k$ (see Eq. (19) in Ref. [6] for details). As claimed in Refs. [17, 18, 19], the constraint on the spectral index is somewhat relaxed, and especially, the spectrum with $n_s = 1$ is marginally inside the 95\% C.L. region.

\footnote{\textit{For the model with }$n = 3$, \textit{the density fluctuations do not determine the inflation scale }$v$, \textit{but it determines the coupling constant }$g$ \textit{as }$g \sim 10^{-(6-7)}$.}
Figure 1: The $k$ dependence of the spectral index $n_s$ for $n = 3 - 8$ and $N_e = 50$. The horizontal grid lines correspond to the result of WMAP three year data [10]. For $k \gtrsim 10^{-2}$, $n_s \simeq 1 - 2k$, and for $k = 0$, $n_s \simeq (N_e(n - 2) - (n - 1))/(N_e(n - 2) + (n - 1))$.

which is measured [10] as

$$\delta = \frac{1}{5\sqrt{3\pi}} |V'| \simeq 1.9 \times 10^{-5}. \quad (11)$$

The left panel of Fig. 2 shows the $g$ dependence of the inflation scale $v$ in Eq. (10) for $n = 4 - 8$ and a given e-folding number $N_e = 50$. We see from the figure that the inflation scale increases as the coupling constant $g$ decreases. We also see that the inflation scale becomes higher for larger $n$.

The e-folding number $N_e$ is related to the inflation scale $v$ and the reheating temperature $T_R$ as

$$N_e \simeq 67 + \frac{1}{3} \ln \frac{v^2}{\sqrt{3}} + \frac{1}{3} \ln T_R. \quad (12)$$

Thus, we can also obtain the inflation scale $v$ for a given $T_R$, instead of for a given $N_e$, by re-solving Eq. (11), although the resultant $T_R$ dependence of $v$ is very weak.

For a given inflation scale $v$ and $n$, we find the gravitino mass from Eq. (7). The right panel of Fig. 2 shows the $g$ dependence of the gravitino mass for $n = 4 - 8$. Similarly to the inflation scale $v$, the gravitino mass also increases as the $g$ decreases and as $n$ increases. Especially, the gravitino mass is roughly given by

$$m_{3/2} \sim 300 g^{-3/2} \text{GeV}, \quad (13)$$

for $n = 4$, while $m_{3/2} \gtrsim 10^6 \text{GeV}$ for $n \geq 5$ and $g < \mathcal{O}(1)$. In the following, we fix $n = 4$, since we are interested in TeV-scale SUSY breaking.
Figure 2: Left: The $g$ dependence of the inflation scale $v$ for $n = 4 - 8$, $N_e = 50$ and $k = 10^{-2}$. The dependence of $v$ on $N_e$ and $k$ are very small. Right: The $g$ dependence of the gravitino mass for $n = 4 - 8$, $N_e = 50$ and $k = 10^{-2}$.

Interestingly, as we see from Eq. (13), the new inflation model with $n = 4$ easily accommodates anomaly-mediation models which are realized for $m_{3/2} \gtrsim 30$ TeV, by taking $g = \mathcal{O}(10^{-2})$. Furthermore, as we will show in the next section, the inflaton decay may provide a sufficient amount of winos which explains the observed dark matter density in the universe.

Let us now discuss a reheating process of the present new inflation. We introduce the following superpotential interaction between the inflaton and the right-handed neutrinos $N$'s,

$$\delta W = \frac{h}{6} \phi^3 N^2,$$

(14)

where $h$ is a dimensionless parameter and we take it positive [6]. Notice that we have introduced another parameter $h$ in the model. At the vacuum, this term Eq. (14) induces masses of the right handed neutrinos as

$$m_N = \frac{h}{3} \phi_0^3 \approx \frac{h}{12g} m_\phi.$$  

(15)

Then, if $2m_N < m_\phi$ (i.e., $h < 6g$) the inflaton decays into a pair of right-handed neutrinos and the reheating occurs after the inflation. The decay rate is given by

$$\Gamma_N \approx \frac{|h|^2}{16\pi} \phi_0^4 m_\phi.$$  

(16)

Consequently, the reheating temperature becomes

$$T_R \approx \left( \frac{10}{g_* \pi^2 \Gamma_N^2} \right)^{1/4} \approx 1.5 \times 10^6 h g^{-5/4} \text{GeV},$$

(17)
where $g_\star (\simeq 228.75)$ is the effective number of massless degrees of freedom, and we have used Eqs. (4), (10) and (16).

A nice point of this reheating process is that the production of right-handed neutrinos via the inflaton decay causes the leptogenesis [20] which results in the baryon asymmetry in the universe. As investigated in Ref. [21], the baryon asymmetry per entropy density $s$ is given by,

$$\eta = \frac{n_B}{s} \simeq 8.2 \times 10^{-11} \left( \frac{T_R}{10^6 \text{GeV}} \right) \left( \frac{2m_N}{m_\varphi} \right) \left( \frac{m_{\nu_3}}{0.05 \text{eV}} \right) \delta_{\text{eff}}. \quad (18)$$

Here $m_{\nu_3}$ is the mass of the heaviest (active) neutrino, the phase $\delta_{\text{eff}} = O(1)$ is the effective CP-violating phase defined in Ref. [20], and we have assumed the ratio of the vacuum expectation value of up- and down-type Higgs bosons to be much larger than 1. Since $T_R$, $m_\varphi$ and $m_N$ are given by $g$ and $h$, the baryon asymmetry can be expressed in terms of $g$ and $h$, or equivalently $m_{3/2}$ and $T_R$, neglecting the weak dependence on $k$.

Fig. 3 shows the allowed parameter region on $(m_{3/2}, T_R)$ plane. The dashed (blue) line is the upper bound on the reheating temperature which comes from the condition,

$$m_N < \frac{m_\varphi}{2}, \quad (i.e., \ h < 6g). \quad (19)$$

The short dashed (green) line denotes the lower bound on the reheating temperature to explain the observed baryon asymmetry $\eta = (8.7 \pm 0.3) \times 10^{-11}$ [10]. As we see from the figure, the allowed region of the reheating temperature is rather constrained,

$$T_R \simeq 10^{6-7} \text{ GeV}. \quad (20)$$

Before closing this section, we summarize the relation between model parameters and the physical quantities. As we see in Eqs. (7), (9), (10) and (17), all the parameters in the model, $k$, $v$, $g$, $n$ and $h$, are determined by the spectral index, the density fluctuations, the gravitino mass and the reheating temperature (equivalently the observed baryon asymmetry). As we have seen so far, four out of the above five parameters are already constrained by observations. Therefore, there remains no free parameter other than the gravitino mass which is one of the most important parameter of the low energy SUSY models and will be determined by future experiments.
Figure 3: The gravitino mass dependence of the upper and lower bound on the reheating temperature. The dashed (blue) line corresponds to the upper bound on the reheating temperature and the short dashed (green) line to the lower bound on the reheating temperature.

3 Dark matter density in anomaly-mediated SUSY breaking

As mentioned in the introduction we adopt anomaly-mediated SUSY breaking where gauginos acquire SUSY-breaking masses through anomalies of the scale invariance \[22, 23\], while squarks and sleptons receive SUSY breaking soft masses directly from SUGRA effects. Thus, there are mass gaps between gauginos and squarks/sleptons \[23\].

In anomaly-mediation models, the most probable lightest supersymmetry particle (LSP) candidate is the neutral wino\(^3\) which has a mass \(m_{\tilde{w}} \simeq 3 \times 10^{-3} m_{3/2} \) and it is a good candidate for the dark matter \[23, 24, 25, 26\].\(^4\) However, the thermally produced winos cannot explain the observed dark matter density unless it is very heavy, \(m_{\tilde{w}} \simeq 2 \text{ TeV} \), because of its large annihilation cross section. The number density of winos produced via the decay of the gravitinos depends on the relic number density of gravitinos. We easily find that the gravitinos produced by scattering processes in the thermal bath cannot supply the observed dark matter density unless \(T_R \gtrsim 10^8 \text{ GeV} \). Thus, if we consider the wino dark matter with \(T_R \simeq 10^{6-7} \text{ GeV} \), the mass of the wino is too heavy \((m_{\tilde{w}} \simeq 2 \text{ TeV})\) to be found in the next generation of collider experiments. Fortunately, in

\(^3\)The finite one-loop corrections to the gaugino masses from the Higgs and Higgsino exchanges can be non-negligible when the supersymmetric Higgs mass \(\mu\) is of the same order of \(m_{3/2} \) \[23\]. In the following, we neglect such one-loop corrections.

\(^4\)The following discussion is almost independent of masses of sfermions as long as the LSP is the neutral wino.
our case, there is another source of the winos, i.e., the gravitinos produced directly from the inflaton decay.

To show it explicitly, we derive the relic density of gravitinos produced directly from the inflaton decay. The relevant terms of the inflaton decay into a pair of gravitinos are,

$$K = |S|^2 + |\phi|^2 + b|S|^2|\phi|^2 + \cdots,$$

(21)

where $S$ is the hidden sector field which has a non-vanishing $F$-term, $b$ a real constant, and the ellipses the higher dimensional terms. Since we have no symmetries to suppress $b$, we naively expect it to be of order one.

As discussed in Ref. [27] the hidden sector fields and the inflaton $\phi$ mix each other and can be made diagonal by the transformation,

$$\hat{\phi} \simeq \phi + \epsilon S, \quad \hat{S} \simeq S - \epsilon^* \phi,$$

(22)

with a mixing angle $\epsilon$,

$$\epsilon \simeq \sqrt{3}(1 + b)\phi_0 \frac{m_{3/2} m_{\phi}}{m_S^2}.$$

(23)

Here, we have assumed the vacuum expectation value of $S$ and the holomorphic mixing mass terms are negligible. Note that we have also assumed that mass of the hidden field, $m_S$, is much larger than the inflaton mass $m_\phi \simeq 10^{10}$ GeV. This assumption is reasonable for the dynamical SUSY breaking models where the hidden sector fields have masses of order of the SUSY breaking scale $\Lambda_{\text{SUSY}} \simeq 10^{12}$ GeV.

The mixing between hidden and inflaton fields leads to an effective coupling of the inflaton to the gravitinos [27],

$$|G^{(\text{eff})}_\phi| \simeq \sqrt{3} \frac{m_S^2}{m_\phi^2} |\epsilon| \simeq 3 \times \frac{m_{3/2}}{m_\phi}(1 + b)\phi_0,$$

(24)

which induces the decay of $\hat{\phi}$ into a pair of gravitinos. The decay rate is given by

$$\Gamma_{3/2} \simeq \frac{|G^{(\text{eff})}_\phi|^2}{288\pi} \frac{m_\phi^5}{m_{3/2}^2 M_G^2}.$$

(25)

Then, the gravitino-entropy ratio (yield) is given by [11, 12],

$$Y_{3/2}^{(\text{inf})} = 2 \frac{\Gamma_{3/2} 3T_R}{\Gamma_N 4m_\phi} \simeq 4.5 \times |G^{(\text{eff})}_\phi|^2 \left(\frac{m_\phi}{10^9 \text{ GeV}}\right)^4 \left(\frac{10^7 \text{ GeV}}{T_R}\right) \left(\frac{\text{TeV}}{m_{3/2}}\right)^2.$$

(26)
By substituting Eq. (24) into Eq. (26), we obtain the yield as,

\[ Y^{(\text{inf})}_{3/2} \simeq \frac{7 \times 10^{-16}}{(1 + b)^2} \left( \frac{m_\phi}{10^9 \text{GeV}} \right)^2 \left( \frac{10^7 \text{GeV}}{T_R} \right) \left( \frac{\phi_0}{10^{16} \text{GeV}} \right)^2, \]  

(27)

\[ \simeq 4 \times 10^{-13} \times (1 + b)^2 \left( \frac{m_{3/2}}{100 \text{ TeV}} \right)^{4/3} \left( \frac{10^7 \text{ GeV}}{T_R} \right), \]  

(28)

where we have also used Eqs. (5), (4), (10) and (13). We see that there are no mass suppression factors discussed in Ref. [27, 28], since the mass of the hidden sector field \( S \) is larger than that of the inflaton. Notice that we have a substantial amount of gravitinos even for the minimal Kähler potential, i.e., for \( b = 0 \).

The above gravitinos decay to the winos, and the resultant yield of the winos is given by\(^5\)

\[ Y_\tilde{w} \simeq Y^{(\text{inf})}_{3/2}. \]  

(29)

There are other contributions to the yield of the winos, one from the thermally produced winos,

\[ Y_\tilde{w} \simeq 10^{-14} \left( \frac{m_\tilde{w}}{100 \text{ GeV}} \right), \]  

(30)

and the other from the decay of gravitinos produced by the thermal scattering [29, 30],

\[ Y_\tilde{w} \simeq 2 \times 10^{-15} \left( \frac{T_R}{10^7 \text{ GeV}} \right). \]  

(31)

By comparing Eqs. (29)-(31), we find that the dominant source of the winos is gravitinos produced directly from the inflaton decay for \( T_R \simeq 10^{6-7} \text{ GeV} \). Then, we find that the mass density parameter of the wino is given by

\[ \Omega_{\tilde{w}} h^2 = \frac{m_\tilde{w} Y_\tilde{w}}{3.5 \times 10^{-9} \text{ GeV}}, \]  

(32)

\[ \simeq 0.04 \times (1 + b)^2 \left( \frac{m_{3/2}}{100 \text{ TeV}} \right)^{7/3} \left( \frac{10^7 \text{ GeV}}{T_R} \right), \]  

(33)

for \( T_R \simeq 10^{6-7} \text{ GeV} \).

Fig. 4 shows the mass density of the wino for \( b = 0 \). We plot the wino mass density in a \((m_{3/2}, T_R)\) plane in the left panel, and the \( T_R \) dependence of the wino density for a

\(^5\)The annihilation process of the non-thermally produced winos is ineffective, since the gravitino decay occurs at very low temperatures, \( T \simeq \mathcal{O}(10) \text{ MeV} \).
given gravitino mass in the right panel. In the figures, we have also scanned $k$ from $10^{-1.5}$ to $10^{-4}$, although the resultant $k$ dependence is negligibly small.

The solid (red) line in the left panel corresponds to the observed dark matter density $\Omega_{\text{DM}} h^2 \simeq 0.127$, the right side of the line to the region of too much dark matter $\Omega_{\tilde{w}} h^2 > 0.127$, and the left side of the line to the region of less dark matter $\Omega_{\tilde{w}} h^2 < 0.127$. For comparison, we also plot the line where $\Omega_{\tilde{w}} h^2 \simeq 0.127$ is satisfied without the gravitinos directly from the inflaton decay as a long-dashed (red) line in the left panel. This line corresponds to the result in Ref. [26], although they have used $m_{\tilde{w}} \simeq 5.2 \times 10^{-3} m_{3/2}$ because of the existence of the light charged particles. From the figures, we find that the mass density of the wino is consistent with the observed dark matter density in a narrow range of the wino mass, $400 \text{GeV} \lesssim m_{\tilde{w}} \lesssim 750 \text{GeV}$, for $T_R \simeq 10^{6-7} \text{GeV}$ and $b = 0$.

As a result, we find that the gravitinos directly from the inflaton decay provide a sufficient amount of the winos for the dark matter if

$$(1 + b)^{-6/7} \times 400 \text{GeV} \lesssim m_{\tilde{w}} \lesssim \text{Min} \left[ (1 + b)^{-6/7} \times 750 \text{GeV}, \ 2.1 \text{TeV} \right].$$

Here, the upper bound on the wino mass on the right hand side, $m_{\tilde{w}} \simeq 2.1 \text{TeV}$, corresponds to the long-dashed (red) line in Fig. 4 where the observed dark matter density is supplied by the thermally produced winos. Since the parameter $b$ is of order one, we consider that the $|b|$ ranges from $1/3$ to $3$. In that case, the mass density of the wino is consistent with the observed dark matter in a wide range of the wino mass, $100 \text{GeV} \lesssim m_{\tilde{w}} \lesssim 2 \text{TeV}$.

4 Conclusions

In this paper, we study the new inflation model [7, 8] which is well consistent with the WMAP observations. The remarkable feature of this model is that there remains essentially only one free parameter, the gravitino mass, and the other parameters are determined by the observations. However, as discussed in Refs. [14], gravity-mediation models for SUSY breaking suffer from a serious gravitino-overproduction problem, and hence, the new inflation model suggests gauge-mediation or anomaly-mediation models.

As to gauge-mediation models, however, they do not well accord with the new inflation model compared to the anomaly-mediation models. Firstly, the gravitino mass in this new
inflation model is rather large, i.e., $m_{3/2} \geq \mathcal{O}(10)\,\text{GeV}$, where the equality is saturated for a large coupling constant $g \simeq 10$ (see Eq. (13)). Secondly, the gravitino is the LSP, but it is difficult to explain the observed dark matter density, since the gravitino abundance is in short supply for $T_R \simeq 10^{6-7}\,\text{GeV}$ and $m_{3/2} \simeq \mathcal{O}(10)\,\text{GeV}$ (see Eq. (27)).

On the contrary, anomaly-mediation models are in harmony with the new inflation model. As we have shown, the new inflation model with $n = 4$ can easily realize the gravitino mass $m_{3/2} \gtrsim 30\,\text{TeV}$ which is suitable for the anomaly-mediation. Furthermore, we have found that gravitinos produced directly from the inflaton decay provide a sufficient amount of the winos for the dark matter in the universe. As a result, we found that the relic density of the wino LSP is consistent with the observed dark matter density in a wide range of the wino mass, $100\,\text{GeV} \lesssim m_{\tilde{w}} \lesssim 2\,\text{TeV}$.

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6 The coupling constant $g$ should be at most $g \lesssim 10$, since otherwise the large coupling constant leads to large radiative corrections to the Kähler potential which invalidates the flatness of the inflaton potential.
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