Enhanced Higgs Mass in a Gaugino Mediation Model without the Polonyi Problem

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Abstract

We consider a SUSY breaking scenario without the Polonyi problem. To solve the problem, the enhanced couplings of the Polonyi field to an inflaton, gauge kinetic functions and itself are assumed. As a result, a gaugino mediated SUSY breaking occurs. In this scenario, the Higgs boson mass becomes consistent with the recently observed value of the Higgs-like boson (i.e., $m_h \simeq 125$ GeV) for the gluino mass about 4 TeV, which is, however, out of the reach of the LHC experiment. We show that the trilinear coupling of the scalar top is automatically enhanced by the presence of the extra matters. With such extra matters, the Higgs mass as large as 125 GeV can be realized with the gluino mass of 1 − 2 TeV which is within the reach of the LHC experiment. In our scenario, the gravitino is the lightest SUSY particle and the candidate for dark matter, and the Wino, Bino, and sleptons are in a range from 200 GeV to 700 GeV.
1 Introduction

The presence of so called Polonyi field $Z$ is an inevitable ingredient in the gravity mediation when the gravitino mass is $\lesssim \mathcal{O}(1) \text{ TeV}$, otherwise we have vanishing gaugino masses at the tree level and the contributions from the anomaly mediation are too small \[1, 2\]. However, this Polonyi field $Z$ causes serious cosmological problems. In particular, its decay in the early universe produces too much entropy, resulting in a huge dilution of the primordial baryon-number asymmetry. Furthermore, its decay occurs during/after the Big-Bang Nucleosynthesis (BBN) and destroys light elements produced by the BBN \[7\].

From a cosmological view point, the Polonyi problem severely restricts the supersymmetry (SUSY) breaking scenarios.

In a series of recent works, it has been pointed out \[8, 9\] that the serious Polonyi problem can be solved if the Polonyi field has enhanced couplings to inflaton. This observation is based on the adiabatic evolution of the Polonyi field following the inflaton potential, originally suggested by Linde \[10\]. The integration of inflaton field induces also enhanced self couplings of the Polonyi field; the Polonyi mass becomes much heavier than the gravitino mass. In addition, enhanced couplings of the Polonyi field $Z$ to the gauge kinetic function are preferred in order to solve the Polonyi problem. This is because the enhanced couplings make the decay of $Z$ faster and the cosmological constraint becomes weaker \[11\]. Consequently, relatively high reheating temperature $T_R \sim 10^6 \text{ GeV}$ is allowed \[9\]. With this reheating temperature, non-thermal leptogenesis works for baryogenesis \[12\].

Such a setup suggests a gaugino mediated SUSY breaking scenario \[13\] (See Refs. \[15\] for a scenario in the framework of extra dimension.) Here, we consider the case where the interactions between the gauge multiplets and the Polonyi field are enhanced while those between chiral multiplets in the minimal supersymmetric standard model (MSSM) and the Polonyi field are not. We study its phenomenological consequences paying par-

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1 In scenarios with $\mathcal{O}(10-100) \text{ TeV}$ gravitino, the Polonyi field is not required since the gaugino masses of $\mathcal{O}(1) \text{ TeV}$ are induced at one-loop level. Without assuming a particular form of the Kahler potential, "Pure Gravity Mediation" scenario \[3, 4\] was proposed where the scalar masses are $\mathcal{O}(10-100) \text{ TeV}$ and the gaugino masses are $\mathcal{O}(1) \text{ TeV}$. (A similar model was proposed in Ref. \[5\].) The scenario can naturally explain the Higgs mass of around 125 GeV with heavy stops \[6\].

2 If we assume a separation between the Polonyi field $Z$ and quark, lepton and Higgs multiplets in the conformal frame, we have vanishing soft masses for sfermions \[13\]. Based on this observation, the gaugino mediation scenario was first proposed in \[13\] (see also \[14\]).
ticular attention to the Higgs mass. As well as the cosmological advantages mentioned above, such a scenario is also favored because it can solve the serious SUSY FCNC problem \cite{13}. As we discuss below, in our setup, the gravitino mass $m_{3/2}$ is much smaller than the gaugino and the Polonyi masses because of the enhanced couplings. The sfermions masses and the scalar trilinear couplings (so-called $A$-terms) are also of the order of the gravitino mass at the tree-level and dominantly arise from the gaugino masses through the renormalization group evolution. Therefore, the SUSY CP and flavor problems are relaxed.

In this letter, motivated by the recent discovery of the Higgs-like boson at the ATLAS and CMS experiments \cite{17}, we investigate the Higgs boson mass in the gaugino mediated SUSY breaking scenario. If the particle content of the MSSM is assumed up to the cut-off scale (which will be taken to be at the GUT scale), the Higgs mass of around 125 GeV is realized only in the region out of the reach of the LHC experiment \cite{18}. However, we point out that the existence of the extra matters can drastically enhance the trilinear coupling of the stop, even if there are no direct couplings to the MSSM matter fields. As a result, the Higgs mass can be as large as 125 GeV with large $A$-term \cite{19} in a region within the reach of the LHC experiment; gluino mass is around 1–2 TeV. The existence of the extra matters at the scale much lower than the Planck scale is expected in many models, for example, like $E_6$ grand unified theory, axion models, and so on.

This paper is organized as follows. In section 2, we explain the setup of our “gaugino dominated SUSY breaking” scenario and show that the Higgs mass of 125 GeV can be consistent only with the gluino mass heavier than about 4 TeV with the particle content of the MSSM \cite{18}. In section 3, we show that the trilinear coupling of the stop is enhanced if the extra matters exist; the Higgs mass can be explained with the gluino mass of 1–2 TeV. The final section is devoted to conclusion and discussion.

\footnote{A similar setup was considered in \cite{16}.}

\footnote{If we assume the sequestered form of the Kahler potential, sfermion masses and $A$-terms vanish at the tree level \cite{13, 11}.}
2 Gaugino mediation without Polonyi problem

Let us discuss the setup of our gaugino mediation scenario. We assume that the Polonyi field $Z$ strongly couples to the inflaton field and the gauge kinetic functions, while the couplings of $Z$ to the other fields (matters and Higgs) are suppressed. The relevant part of the Kahler potential is given by

$$K = -c_I^2 |Z|^2 |I|^2 M_P^2 - c_Z^2 |Z|^4 4M_P^2,$$

(1)

where $Z$ and $I$ are the Polonyi field and the inflaton field, respectively, and $M_P \simeq 2.4 \times 10^{18}$ GeV is the reduced Plank scale. The coefficient $c_I$ is required to be as large as $\sim 100$ so that the Polonyi abundance is sufficiently suppressed by the adiabatic suppression mechanism [10, 8, 9]. The second term arises by the radiative corrections from the inflaton loops and $c_Z$ is also expected to be $\sim 100$. As a result, the Polonyi mass is enhanced compared to the gravitino mass:

$$m_Z = c_Z F_Z M_P \sim c_Z \sqrt{3} m_{3/2},$$

(2)

where $m_{3/2}$ is the gravitino mass. Here we assume that the SUSY is dominantly broken by the $Z$ field so that $m_{3/2} \simeq F_Z / \sqrt{3} M_P$. The couplings of $Z$ to the gauge kinetic functions are also assumed to be enhanced:

$$\mathcal{L} \ni -c_g \int d^2 \theta Z W^a W^a M_P + h.c.,$$

(3)

where $c_g \sim 100$. (Here we take the basis in which the gauge kinetic function is canonically normalized.) With Eq. (3), the gaugino mass is given by

$$M_\lambda = 2 c_g F_Z M_P \sim 2 \sqrt{3} c_g m_{3/2}.$$

(4)

Since we consider the scenario that $c_g \sim c_Z \sim 100$, the gravitino is the LSP and candidate for a dark matter. The Polonyi field decays into SM gauge bosons through the operator (3), and its decay width is given by

$$\Gamma(Z \rightarrow 2 \text{ gauge bosons}) \simeq c_g^2 \frac{3 m_Z^3}{2 \pi M_P^2} \simeq 1.2 \text{ sec}^{-1} \left( \frac{c_g}{100} \right)^2 \left( \frac{m_Z}{10^3 \text{ GeV}} \right)^3.$$

(5)
With the suppression of the Polonyi abundance and the relatively short lifetime (less than 1 second), the BBN constraint can be avoided relatively easily \cite{11} and the reheating temperature $T_R \simeq 10^6 \text{GeV}$ is allowed for $m_{3/2} \gtrsim 30 \text{GeV}$ and $c_1 \simeq c_Z \simeq c_g \simeq 100$ \cite{9}. With such a relatively high reheating temperature, enough baryon asymmetry can be generated by the non-thermal leptogenesis \cite{12}.

In our setup, the gravitino mass as well as the scalar masses are suppressed compared to the gaugino masses at the cut-off scale. Then the scalar masses dominantly arise from the renormalization group effect between the cut-off scale and the SUSY scale (which is the mass scale of the MSSM SUSY particles) of $\mathcal{O}(1) \text{TeV}$. We calculate the low-energy SUSY parameters by numerically solving two-loop renormalization group equations (RGEs). Here, we use \texttt{SuSpect} \cite{20} to evaluate the spectrum of the SUSY particles. The boundary condition is taken such that all the scalar masses and trilinear coupling constants vanish and the gaugino masses are universal at the cut-off scale (which is taken to be the GUT scale). Then, we calculate the lightest Higgs boson mass using \texttt{FeynHiggs} \cite{21}. In fig. 1, the contours of the constant Higgs mass are shown. In the same figure, we also plot the contours of constant $B$-parameter at the GUT scale, where the $B$-parameter is defined as

\begin{equation}
V \ni B \mu H_u H_d + h.c.,
\end{equation}

with $H_u$ ($H_d$) being the up-type (down-type) Higgs and $\mu$ being the Higgsino mass parameter in the superpotential. (Here, $\mu$ is assumed to be a free parameter, and is added by hand.) The GUT-scale value of the $B$-parameter (which is denoted as $B_{\text{GUT}}$) is expected to be of the same order of the gravitino mass. As we have mentioned, $m_{3/2} \gtrsim 30 \text{GeV}$ is required for non-thermal leptogenesis with avoiding the BBN constraints. So, we take $B_{\text{GUT}} = \pm 30 \text{GeV}$ as representative values.

Adapting the uncertainty in the Higgs mass calculation of $2 - 3 \text{ GeV}$ \cite{22, 23}, Higgs mass as large as $m_h \simeq 125 \text{ GeV}$ may be realized if the gluino mass is heavier than about 4 TeV \cite{18}. Unfortunately, such a heavy gluino (and squarks) is out of the reach of the LHC experiments.
Enhanced $A$-term from extra matters

So far, we have seen that, if we adopt the particle content of the MSSM, the Higgs mass of $m_h \simeq 125$ GeV is hardly realized in gaugino mediation model with gluino and squarks which are within the reach of the LHC experiment. Now, we show that such a conclusion is altered if there exist extra vector-like multiplets at around $1-10$ TeV. Many models such as $E_6$ grand unified theory and axion models predict the existence of the extra matters at the scale much lower than the Planck scale.

The existence of extra matters may enhance the Higgs mass via two effects. First, if the extra matters have sizable Yukawa interaction with the Higgs fields, the radiative correction below the SUSY scale may significantly enhance the lightest Higgs mass [24]. This is the case with $10 + \overline{10}$ extra matters, for example. Second, the presence of the vector-like multiplets changes the beta-functions of the gauge couplings and the gauginos. Consequently, the trilinear coupling of the stop becomes larger than the case without extra matters, resulting in the enhancement of the Higgs mass. The first effect has been intensively studied in recent works [24], so we concentrate on the second one, assuming that the extra matters have no Yukawa interaction with the Higgs fields.
First, we show how $A$-parameters, a squark mass and the ratio of the $A$-parameter to the squark mass $A_q/m_q$ are enhanced with extra matters. For this purpose, we use one-loop RGEs (although our numerical calculations are performed at the two-loop level). With the presence of the extra matters, the beta-functions of the gauge coupling constants and gaugino masses are given by

$$\frac{dg_a^2}{d\ln Q} = (b_a + N_5) \frac{g_a^4}{8\pi^2},$$  

(7)  

$$\frac{dM_a}{d\ln Q} = (b_a + N_5) \frac{g_a^2}{8\pi^2} M_a,$$  

(8)  

where $g_a$ ($a = 1, 2, 3$) are the gauge couplings for $U(1)_Y$, $SU(2)_L$ and $SU(3)_C$, respectively, and $b_a = (33/5, 1, -3)$ are the coefficients of the beta-functions with the MSSM matter content. The number of the extra vector-like multiplets in units of fundamental and anti-fundamental representation of $SU(5)$ GUT gauge group, $5 + \overline{5}$, is denoted as $N_5$. For $N_5 \gtrsim 3$, $g_3$ and $M_3$ at the SUSY scale are smaller than those at the GUT scale. The change of these beta-functions dramatically alters the squark masses and trilinear couplings at the SUSY scale. In particular, the changes of the stop masses and stop trilinear coupling lead to the important consequences in the Higgs boson mass and the SUSY search at the LHC.

Neglecting Yukawa couplings, the RGE of an $A$-parameter of a squark can be written as

$$\frac{dA_q}{d\ln Q} = \frac{1}{16\pi^2} \left( c_a g_a^2 M_a \right)$$

$$= \frac{1}{16\pi^2} \left[ c_a \left( \frac{8\pi^2}{b_a + N_5} \right) \frac{dg_a^2}{d\ln Q} \left( \frac{M_a}{g_a^2} \right) \right].$$  

(9)  

Notice that $M_a/g_a^2$ is an RGE invariant quantity, i.e., constant. The coefficient $c_3 = 32/3$ is common to all $A_q$. (Here, we neglect the effects of Yukawa coupling constants, which do not change the following discussion qualitatively. Our numerical calculation will be performed with the effects of Yukawa coupling constants.) Eq. (9) can be solved as

$$A(Q) = -\frac{c_a}{2(b_a + N_5)} \left[ g_a^2(M_{\text{GUT}}) - g_a^2(Q) \right] \left( \frac{M_a}{g_a^2} \right).$$  

\footnote{Although the one-loop beta-function of $g_3$ vanishes for $N_5 = 3$, inclusion of the higher order corrections leads to the positive beta-function.}
Neglecting the contributions from fixed gluino mass.

A of the gluino mass at the SUSY scale. The ratio

to be 1.2 TeV and $\tan \beta$ increases, because of the larger ratio of $A$. Consequently, the Higgs boson mass is enhanced as the number of the extra matters

takes the renormalization scale $Q$ as the mass of the extra matter, $M_{N_5}$, i.e., the decoupling scale, larger $N_5$ results in a larger $A_q(M_{N_5})$ for the fixed value of $(M_a/g_a^2)$, since $g_a(M_{N_5})$ does not depend on $N_5$; the trilinear couplings including $A_t$ are enhanced for the fixed gluino mass.

Similarly, the RGE of a squark mass can be written as

$$\frac{dm^2_q}{d\ln Q} = \frac{1}{16\pi^2} \left(-d_a g_a^2 M_a^2\right)$$

where $d_3 = 32/3$ is common to all squark mass. By solving Eq. (11), we obtain

$$m^2_q(Q) = \frac{1}{4} \left(\frac{d_a}{b_a + N_5}\right) \left[ g_a^4(M_{GUT}) - g_a^4(Q) \right] \left(\frac{M_a}{g_a^2}\right)^2$$

Again, $m^2_q(M_{N_5})$ becomes larger as $N_5$ increases for the fixed value of $M_a/g_a^2$. Thus by adding extra matters, the stop mass $m_{\tilde{t}}$ is expected to be larger with the fixed value of the gluino mass at the SUSY scale. The ratio $A_q(M_{N_5})/m_{\tilde{q}}(M_{N_5})$ is also enhanced. Neglecting the contributions from $g_1$ and $g_2$, the ratio $(A_q(Q)/m_{\tilde{q}}(Q))^2$ becomes

$$\left(\frac{A(Q)}{m_{\tilde{q}}(Q)}\right)^2 = \frac{c_3^2}{d_3(b_3 + N_5)} \left(1 - g_3^2(Q)/g_3^2(M_{GUT})\right) \left(1 + g_3^2(Q)/g_3^2(M_{GUT})\right)^{-1}$$

Taking $Q = M_{N_5}$, we obtain the enhanced ratio $[A_q(M_{N_5})/m_{\tilde{q}}(M_{N_5})]^2$ for larger $N_5$. Consequently, the Higgs boson mass is enhanced as the number of the extra matters increases, because of the larger ratio of $A_t/m_{\tilde{t}}$ and the larger $m_{\tilde{t}}$.

The results of the numerical calculations are shown in fig. 2. The gluino mass is fixed to be 1.2 TeV and $\tan \beta = 25$ in both left and right panels. In the left panel, the Higgs
mass as a function of $M_{N_5}$ is shown. Three curves correspond to $N_5 = 3, 4, 5$ from bottom to top. For comparison, the dashed line, which is evaluated in the MSSM is also drawn. Remarkably, the Higgs mass reaches 125 GeV with $N_5 = 4$ and $M_{N_5} \simeq 3$ TeV. In the right panel, we also show the normalized trilinear coupling $X_t/m_{\tilde{t}}$, where $m_{\tilde{t}} \equiv (m_{\tilde{t}_1} + m_{\tilde{t}_2})/2$ and $X_t = A_t - \mu \cot \beta$, with $m_{\tilde{t}_1}$ and $m_{\tilde{t}_2}$ being the lighter and heavier stop masses. Three curves correspond to $N_5 = 3, 4, 5$ from top to bottom. The enhancement of $X_t/m_{\tilde{t}}$ can be seen. The number of the vector-like multiplets $N_5 = 3$ is motivated by the $E_6$ grand unified theory, and $N_5 = 4$ can be regarded as a complete vector-like family.

The contours of the constant Higgs mass for $N_5 = 3$ and $M_{N_5} = 1$ TeV are shown in fig. 3 (left panel). The regions near the lines $B_{GUT} = \pm 30$ GeV are consistent with the boundary condition of the gaugino mediated SUSY breaking scenario. The Higgs mass is calculated to be $123 - 124$ GeV for the gluino mass of $1.1 - 1.6$ TeV. Correspondingly, a squark mass is $1.8 - 2.6$ TeV (right panel). Notice that these gluino and squarks may be observed at the LHC. In this region, the mass of the (right-handed) slepton is about $470 - 670$ GeV and the Bino (Wino) mass is $210 - 310$ GeV ($330 - 490$ GeV), provided the universal gaugino masses at the GUT scale. Such relatively light non-colored SUSY particles may be seen at future $e^+ e^-$ linear collider experiments.

We also show that contours of the constant Higgs mass for $N_5 = 4$ and $M_{N_5} = 4$ TeV in fig. 4. The Higgs mass of about 125 GeV (or larger) is realized in a wide region where the gluino mass is larger than about 1.2 TeV. Correspondingly, the Bino and Wino masses are larger than 380 GeV and 630 GeV, respectively. Notice that the squark mass is larger than 3 TeV, which is too heavy to be observed at the LHC. The sleptons are also heavier than about 1 TeV. In the calculation, we adapt the universal gaugino masses at the GUT scale and the Bino is the next-to-the lightest SUSY particle (NLSP) in the whole region. (For the case without the GUT relation, see the discussion below.)

So far, we have seen that the existence of extra matters at the SUSY scale enhances the Higgs mass in the gauge mediated SUSY breaking scenario. More enhancement may be realized if there exist additional extra matters at the intermediate scale. One of the motivations to consider such extra matters at the intermediate scale is SUSY axion model. As an example of the enhancement due to the extra matters at the intermediate scale, we
Figure 2: The Higgs boson mass and the normalized trilinear coupling of the stop as a function of the decoupling scale of the extra matter. The gluino mass is fixed to be $m_{\tilde{g}} = 1.2$ TeV. Here, $\tan \beta = 25$.

consider the case where there exists one pair of $5 + \bar{5}$ extra matters at $10^8$ GeV (as well as three pairs of $5 + \bar{5}$ extra matters at TeV scale). In fig. 5, contours of the Higgs mass (left panel) and squark mass (right panel) are shown.

Finally, we comment on cosmological implication of Bino-like neutralino as the NLSP. In the early universe, neutralinos are produced and may decay after the BBN epoch in particular when R-parity is conserved. Hadro- and photo-dissociation processes are caused by the decay of the neutralino, and the success of the BBN may be spoiled if the lifetime of the neutralino is longer than $\sim 1$ sec \[11\]. In the present case where $M_a/m_{3/2}/2 \sim 100$, the parameter region which we are interested in mostly conflicts with the BBN constraints. Such a problem may be solved if the NLSP is Wino-like neutralino.

\[6\] The mass can be smaller than the Peccei-Quinn (PQ) symmetry breaking scale, i.e., for instance, Yukawa coupling of $O(10^{-2}) \times$ PQ breaking scale.

\[7\] If stau is the NLSP, which may be the case in the gaugino mediation model without extra matters, the constraint is much weaker \[11\]. Assuming that stau decays into tau and gravitino via supercurrent interaction, the stau mass is required to be larger than 200 GeV for $m_{\tilde{\tau}}/m_{3/2} = 100$ (with $m_{\tilde{\tau}}$ being the mass of stau), in order for successful BBN scenario. In addition, for $m_{\tilde{\tau}}/m_{3/2} = 300$, no constraint is obtained.
instead of Bino-like neutralino. In such a case, the thermal relic abundance of the NLSP is significantly suppressed, and the BBN constraints are relaxed. Interestingly, if the Wino-like neutralino is the NLSP, the signal of the Wino production may be observed at the LHC [28, 4]. Another possibility to avoid the confliction with the BBN constraints is to introduce small R-parity violation.

4 Conclusion and discussion

In this letter, we have considered a gaugino mediated SUSY breaking scenario without the Polonyi problem. The gaugino mediation naturally occurs with the requirements for avoiding the serious Polonyi problem. With this setup, we have evaluated the Higgs boson mass and found that the Higgs mass of around 125 GeV may be realized with gluino mass heavier than about 4 TeV in the MSSM. Unfortunately, such heavy colored SUSY particles are out of reach of the LHC experiment.

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8The GUT relation among gaugino masses may not hold even if the SM gauge couplings are unified. For example, in the product-group unification scenario [26], this is the case [27].
Figure 4: Contours of the Higgs mass (left) and the squark mass (right) on $m_{\tilde{g}} - \tan \beta$ plane in the unit of GeV. The four pairs of the extra matters exist at 4 TeV.

However, if there exist a number of extra matters at $1 - 10$ TeV, the Higgs boson mass can be significantly enhanced with relatively small gluino mass of $1 - 2$ TeV. This is because the trilinear coupling of the scalar top is enhanced by the change of the RGEs of the gauge coupling constants and gaugino masses. With the uncertainty in the calculation of the Higgs boson mass ($2 - 3$ GeV), the squark mass of around 2 TeV becomes consistent with observed value of the Higgs boson mass. Such gluino and squarks are within the reach of the LHC experiment. In addition, it is notable that the masses of other SUSY particles can be much below 1 TeV; in the parameter region of our interest, the masses of slepton, Bino, and Wino are about $470 - 670$ GeV, $210 - 310$ GeV, and $330 - 490$ GeV, respectively. These non-colored SUSY particles can be targets of future $e^+e^-$ linear collider experiments.

In this letter, we have concentrated on the scenario in which only the gauge multiplets (and the inflaton) have enhanced couplings to the Polonyi field. From the point of view of solving the SUSY FCNC problem, the Higgs fields may also strongly couple to the Polonyi field. If so, the behaviors of the SUSY breaking parameters may be significantly altered \[18\]. In particular, the cut-off-scale values of the $A$-parameters and soft SUSY
Figure 5: Contours of the Higgs mass (left) and the squark mass (right) on $m_{\tilde{g}} - \tan \beta$ plane in the unit of GeV. The three pairs of the extra matters exist at 1 TeV and another pair of the extra matters exist at $10^8$ GeV.

Breaking Higgs masses are expected to become much larger than the gravitino mass. Such a scenario will be studied in a separate publication [29].

Our scenario is consistent with the cosmological observation. The baryon asymmetry can be from non-thermal leptogenesis with relatively high reheating temperature as $10^6$ GeV. The gravitino is the LSP and the candidate for dark matter. If the R-parity is conserved, the gravitino mass of $m_{3/2} \sim 10$ GeV, which is expected in the present scenario, may conflict with the BBN constraints in particular when the Bino-like neutralino is the NLSP. However, the BBN constraints can be avoided if the NLSP is Wino-like neutralino or if a small violation of R-parity is introduced. Even with a small R-parity violation, the gravitino can be long-lived enough to be a viable candidate for dark matter [30, 31].
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References

[1] L. Randall and R. Sundrum, Nucl. Phys. B 557, 79 (1999) [hep-th/9810155].
[2] G. F. Giudice, M. A. Luty, H. Murayama and R. Rattazzi, JHEP 9812, 027 (1998) [hep-ph/9810442].
[3] M. Ibe and T. T. Yanagida, Phys. Lett. B 709, 374 (2012) [arXiv:1112.2462 [hep-ph]]; M. Ibe, S. Matsumoto and T. T. Yanagida, Phys. Rev. D 85, 095011 (2012) [arXiv:1202.2253 [hep-ph]].
[4] B. Bhattacherjee, B. Feldstein, M. Ibe, S. Matsumoto and T. T. Yanagida, arXiv:1207.5453 [hep-ph].
[5] L. J. Hall, Y. Nomura and S. Shirai, arXiv:1210.2395 [hep-ph].
[6] Y. Okada, M. Yamaguchi and T. Yanagida, Prog. Theor. Phys. 85, 1 (1991); J. R. Ellis, G. Ridolfi and F. Zwirner, Phys. Lett. B 257, 83 (1991); H. E. Haber and R. Hempfling, Phys. Rev. Lett. 66, 1815 (1991).
[7] G. D. Coughlan, W. Fischler, E. W. Kolb, S. Raby, G. G. Ross, Phys. Lett. B131, 59 (1983); J. R. Ellis, D. V. Nanopoulos, M. Quiros, Phys. Lett. B174, 176 (1986); A. S. Goncharov, A. D. Linde, M. I. Vysotsky, Phys. Lett. B147, 279 (1984).
[8] K. Nakayama, F. Takahashi and T. T. Yanagida, Phys. Rev. D 84, 123523 (2011) arXiv:1109.2073 [hep-ph].
[9] K. Nakayama, F. Takahashi and T. T. Yanagida, Phys. Rev. D 86, 043507 (2012) arXiv:1112.0418 [hep-ph].
[10] A. D. Linde, Phys. Rev. D 53, 4129 (1996) [hep-th/9601083].
[11] M. Kawasaki, K. Kohri and T. Moroi, Phys. Lett. B 625, 7 (2005) [astro-ph/0402490]; M. Kawasaki, K. Kohri and T. Moroi, Phys. Rev. D 71, 083502 (2005) [astro-ph/0408426]; M. Kawasaki, K. Kohri, T. Moroi and A. Yotsuyanagi, Phys. Rev. D 78, 065011 (2008) [arXiv:0804.3745 [hep-ph]].

[12] K. Kumekawa, T. Moroi and T. Yanagida, Prog. Theor. Phys. 92, 437 (1994) [hep-ph/9405337]; K. Hamaguchi, H. Murayama and T. Yanagida, Phys. Rev. D 65, 043512 (2002) [hep-ph/0109030].

[13] K. Inoue, M. Kawasaki, M. Yamaguchi and T. Yanagida, Phys. Rev. D 45, 328 (1992).

[14] H. Murayama, H. Suzuki, T. Yanagida and J. i. Yokoyama, Phys. Rev. D 50, 2356 (1994) [hep-ph/9311326].

[15] D. E. Kaplan, G. D. Kribs and M. Schmaltz, Phys. Rev. D 62, 035010 (2000) [hep-ph/9911293]; Z. Chacko, M. A. Luty, A. E. Nelson and E. Ponton, JHEP 0001 (2000) 003 [hep-ph/9911323].

[16] W. Buchmuller, K. Hamaguchi and J. Kersten, Phys. Lett. B 632, 366 (2006) [hep-ph/0506105].

[17] G. Aad et al. [ATLAS Collaboration], Phys. Lett. B 716, 1 (2012); S. Chatrchyan et al. [CMS Collaboration], Phys. Lett. B716, 30 (2012).

[18] F. Brummer, S. Kraml and S. Kulkarni, JHEP 1208, 089 (2012) [arXiv:1204.5977 [hep-ph]].

[19] Y. Okada, M. Yamaguchi and T. Yanagida, Phys. Lett. B 262, 54 (1991).

[20] A. Djouadi, J. -L. Kneur and G. Moultaka, Comput. Phys. Commun. 176, 426 (2007) [hep-ph/0211331].

[21] S. Heinemeyer, W. Hollik and G. Weiglein, Comput. Phys. Commun. 124, 76 (2000) [hep-ph/9812320]. S. Heinemeyer, W. Hollik and G. Weiglein, Eur. Phys. J. C 9, 343 (1999) [hep-ph/9812472]. G. Degrassi, S. Heinemeyer, W. Hollik, P. Slavich and G. Weiglein, Eur. Phys. J. C 28, 133 (2003) [hep-ph/0212020]. M. Frank, T. Hahn, S. Heinemeyer, W. Hollik, H. Rzehak and G. Weiglein, JHEP 0702, 047 (2007) [hep-ph/0611326].
[22] B. C. Allanach, A. Djouadi, J. L. Kneur, W. Porod and P. Slavich, JHEP 0409, 044 (2004) [hep-ph/0406166].

[23] S. P. Martin, Phys. Rev. D 75, 055005 (2007) [hep-ph/0701051].

[24] T. Moroi and Y. Okada, Mod. Phys. Lett. A 7, 187 (1992); Phys. Lett. B 295, 73 (1992); K. S. Babu, I. Gogoladze, M. U. Rehman and Q. Shafi, Phys. Rev. D 78, 055017 (2008) [arXiv:0807.3055 [hep-ph]]; S. P. Martin, Phys. Rev. D 81, 035004 (2010) [arXiv:0910.2732 [hep-ph]].

[25] M. Asano, T. Moroi, R. Sato and T. T. Yanagida, Phys. Lett. B 705, 337 (2011) [arXiv:1108.2402 [hep-ph]]; M. Endo, K. Hamaguchi, S. Iwamoto and N. Yokozaki, Phys. Rev. D 84, 075017 (2011) [arXiv:1108.3071 [hep-ph]]; J. L. Evans, M. Ibe and T. T. Yanagida, [arXiv:1108.3437 [hep-ph]]; T. Moroi, R. Sato and T. T. Yanagida, Phys. Lett. B 709, 218 (2012) [arXiv:1112.3142 [hep-ph]]; M. Endo, K. Hamaguchi, S. Iwamoto and N. Yokozaki, Phys. Rev. D 85, 095012 (2012) [arXiv:1112.5653 [hep-ph]]; S. P. Martin and J. D. Wells, arXiv:1206.2956 [hep-ph].

[26] T. Yanagida, Phys. Lett. B 344, 211 (1995) [hep-ph/9409329].

[27] N. Arkani-Hamed, H.-C. Cheng and T. Moroi, Phys. Lett. B 387, 529 (1996) [hep-ph/9607463].

[28] T. Moroi and K. Nakayama, Phys. Lett. B 710, 159 (2012) [arXiv:1112.3123 [hep-ph]].

[29] T. Moroi, T. T. Yanagida and N. Yokozaki, in preparation.

[30] W. Buchmuller, L. Covi, K. Hamaguchi, A. Ibarra and T. Yanagida, JHEP 0703, 037 (2007) [hep-ph/0702184 [HEP-PH]].

[31] F. Takayama and M. Yamaguchi, Phys. Lett. B 485, 388 (2000) [hep-ph/0005214].