A Note on Polonyi Problem

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Abstract

We reinvestigate the cosmological Polonyi problem in the case where the Polonyi mass is $\mathcal{O}(10)$ TeV. Such a large supersymmetry breaking scale implies that the Polonyi field should be sequestered from the standard model sector. Since the Polonyi field does not have a coupling to the gauge multiplets at tree level, in order to obtain sufficiently high reheating temperature compatible with the standard big-bang nucleosynthesis the Polonyi mass well exceeds 100 – 1000 TeV, depending on the decay channels. Moreover, we find that the branching ratio of the Polonyi field into neutralinos is of order unity, and thus the resulting neutralino LSPs, if stable, overclose the Universe even for the case of the wino-like LSP. Our explicit computation given here exhibits a very serious cosmological difficulty for models where supersymmetry breaking is caused by the Polonyi-type field.
Supersymmetry (SUSY) is one of the most attractive candidates for new physics beyond the standard model (SM), which solves the naturalness problem associated with the electroweak scale. Null results of superparticle searches, however, require it to be broken. A natural thought is that the supersymmetry is broken spontaneously in supergravity (SUGRA) [1], where the super-Higgs mechanism is operative. As a consequence, the would-be Nambu-Goldstone fermion (goldstino) arises, which becomes the longitudinal component of the gravitino. The properties of the scalar superpartner of the goldstino are quite sensitive to how supersymmetry is broken. In particular, when the supergravity effect plays an essential role in supersymmetry breaking (termed the non-renormalizable hidden sector in Ref. [2]), the scalar superpartner of the goldstino often has interaction whose strength is similar to gravitational interaction, a mass, $m_\phi$, of the order of the gravitino mass, $m_{3/2}$, and a vacuum expectation value of the Planck scale. Such a field is called the Polonyi field [3].

In this scenario, the lifetime of the Polonyi field, $\tau_\phi \sim M_{Pl}^2/m_\phi^3$, may become very long, where $M_{Pl} \simeq 2.4 \times 10^{18}$ GeV is the reduced Planck mass. It is well-known that the Polonyi field with the mass of the order of the electroweak scale causes the cosmological difficulty. In inflationary epoch, the Polonyi field is shifted away from the true vacuum with a magnitude, $\phi_{in}$, of the order of the Planck scale, $\phi_{in} \simeq O(M_{Pl})$. After inflation, the Polonyi field starts a coherent oscillation around the true minimum when the Hubble parameter becomes comparable to the Polonyi mass. Its energy density will dominate the energy of the Universe soon after the start of the oscillation. The subsequent decay of the Polonyi field releases tremendous amount of entropy after nucleosynthesis, and thus it jeopardizes the success of the big-bang nucleosynthesis (BBN). This is the notorious Polonyi problem [2] [4] [5].

It is known that the problem is somewhat ameliorated when the Polonyi mass is as heavy as $O(10)$ TeV [6] [7] [8]. In this case, the reheating temperature by the Polonyi decay would be sufficiently high so that its decay would not affect the BBN. The baryon number asymmetry may be diluted at the reheating by the Polonyi decay, but can still be as large as the observed value in the Affleck-Dine mechanism [9] (See, for instance, Ref. [10] [11] for a recent work.). A difficulty in this case is the over-abundance of the lightest superparticles (LSPs) which are produced at the Polonyi decay. The annihilation among them does not work sufficiently, and thus too much amount of the LSPs tends to remain in particular when the LSP is a bino-like neutralino.
Recently, it was recognized that the gravitino production at the decay of the heavy scalar field (the Polonyi field, the modulus field and the inflaton) is sizably large, which is incompatible with the big-bang cosmology \[12\] \[13\] \[14\] \[15\] \[16\] \[17\]. To avoid the gravitino overproduction in the present situation, one should consider the case where the Polonyi decay into the gravitino pair is kinematically forbidden, that is, \(2m_{3/2} > m_\phi\). Assuming that the gravitino mass is a few tens TeV or lighter, this condition implies that the Polonyi field cannot be arbitrarily heavy, but can at most be as large as 100 TeV. On the other hand, with such a large supersymmetry breaking scale, the Polonyi field should be sequestered from the SM sector to keep the superparticle masses not very far from the electroweak scale. Absence of the direct coupling, however, raises the question how fast the Polonyi field can decay into the SM particles.

The purpose of this paper is to investigate the Polonyi decay and discuss its implications to cosmology. The decay of heavy scalar fields, including an inflaton, via supergravity interaction was been investigated in detail in Refs. \[12\] \[13\] \[18\] \[19\] \[20\]. Thus, we apply those results to this particular set-up. Since the Polonyi field does not have a coupling to the gauge multiplets at tree level, it will turn out that, in order to obtain sufficiently high reheating temperature the Polonyi mass well exceeds 100 TeV, even if we consider efficient decay into right-handed neutrinos. In the absence of this decay channel, the Polonyi field dominantly decays into gauge multiplets. In this case, the Polonyi mass should be larger than 1000 TeV. We will also discuss the partial decay into the superpartners in the SM sector. We find that its branching ratio is generically of order unity, and thus the resulting relic abundance of the neutralino LSPs, if stable, tends to be too large, even when the LSP is a wino-like neutralino and its annihilation process is most efficient.

To discuss the effect of the Polonyi field decay on the cosmology, we will consider the couplings of the Polonyi field to the various particles. If there were an unsuppressed coupling between the Polonyi field and the gauge kinetic term, the gaugino mass would be as large as 100 TeV. To avoid it, we assume that the gauge kinetic function does not depend on the Polonyi field \(\phi\). Let us then discuss the coupling of \(\phi\) to chiral multiplets. To obtain TeV scale soft scalar masses, we assume that the visible sector is sequestered from the Polonyi...
field, which implies the following form of the SUGRA $f$ function and the superpotential $W$

$$f = -3 \xi(\phi, \phi^*) + Q_i^\dagger e^{V}Q_i, \quad (1)$$
$$W = W(\phi) + \frac{1}{2} M N N + \frac{1}{6} Y_{ijk} Q^i Q^j Q^k, \quad (2)$$

where $\xi$ is an arbitrary function of $\phi$ and $\phi^*$, and $Q$'s denote matter fields. Here, $g$ and $Y_{ijk}$ are the gauge coupling constant and the Yukawa coupling constant, respectively. We have introduced the right-handed neutrino, $N$, with mass $M$, which will play an essential role in our discussion. The function $f$ is related to the Kähler potential, $K$, via

$$f = -3 e^{-K/3}, \quad (3)$$

where we used the Planck unit $M_{Pl} = 1$. In what follows, we will use this unit unless we explicitly mention. From eq. (3), we can obtain the Kähler potential as

$$K = -3 \ln \xi(\phi, \phi^*) + \frac{1}{\xi} Q_i^\dagger e^{V}Q_i + \cdots. \quad (4)$$

Now, we would like to discuss the decay of the Polonyi field. First, we consider the decay into the matter fermions. Relevant terms for producing the matter fermions are

$$\mathcal{L} = -i g_{ij^*} \bar{\chi}^{\dagger j} \bar{\sigma}^\mu D_\mu \chi^i + \frac{i}{4} g_{ij^*} \left( K_k \partial_\mu \phi^k - K_k \partial_\mu \phi^{* k} \right) \bar{\chi}^j \bar{\sigma}^\mu \chi^j - i g_{ij^*} \Gamma_{k l}^i \left( \partial_\mu \phi^k \right) \bar{\chi}^j \bar{\sigma}^\mu \chi^j$$

$$- \frac{1}{2} e^{K/2} \left( \mathcal{D}_j W \right) \chi^j + h.c., \quad (5)$$

where $D_\mu = \partial_\mu - ig_{ij^*} \Gamma_{k l}^i \left( \partial_\mu \phi^k \right)$ and the covariant derivative $D_\mu = \partial_\mu - ig A^{(a)}_{\mu} T^{(a)}$ with the representation matrix, $T^{(a)}$, for the generator of the gauge group. The sum over indices is understood and $\phi^i$ represent all scalar fields including the Polonyi field. When the Polonyi field $\phi$ is written explicitly, $\phi^i$ and $\chi^i$ represent only the matter fields. From eq. (5) and using the equation of motion, we can obtain interactions of $\phi$ with chiral fermions

$$\mathcal{L} = -\frac{1}{2} e^{K/2} \left( K_{ij} W_{ij} - 2 \Gamma_{ij^*}^{k} W_{jk} \right) \phi \chi^i \chi^j + h.c., \quad (6)$$

where we set $K_{ij}, W_i \ll 1$ because of the assumption that the matter fermions are charged under some symmetries. Eq. (6) shows that the decay amplitude is proportional to the mass of a final-state fermion. Thus the decay into quarks and leptons in the SM sector are
suppressed. By the same reason the three-body decay such as $\phi \to q\bar{q}g$ with $q$ and $g$ being a quark and a gauge boson, respectively, are also suppressed. On the other hand, the decay into the right-handed neutrino pair can be sizable, if it is kinematically allowed. The decay width is

$$\Gamma(\phi \to \nu_R\nu_R) = \frac{\lambda^2}{32\pi} N_f \sqrt{1 - \frac{4M^2}{m^2} \left( 1 - \frac{2M^2}{m^2} \right) \frac{M^2}{m^2} \frac{m^3_\phi}{M_{Pl}^2}},$$  \hspace{1cm} (7)

where we have written the Planck mass explicitly, and $\lambda \equiv \xi_\phi/\xi$ and $N_f$ is the number of the right-handed neutrinos.

On the other hand, interactions of $\phi$ with matter scalars come from the kinetic term and the scalar potential

$$\mathcal{L} = -g_{ij}D_\mu\phi^iD^\mu\phi^j - e^K [g^{ij}(D_iW)(D_jW)^* - 3|W|^2].$$  \hspace{1cm} (8)

We can obtain the interactions for the decay $\phi \to \phi^i\phi^j$ and $\phi \to \phi^i\phi^j g$ from the kinetic term of eq. (8). As in the case of fermions, one can also check that these decay widths are suppressed by the masses of the final-state scalars. Thus, they are also subdominant component of the total decay width of $\phi$ unless the final-state scalars have quite large soft SUSY breaking masses. The interactions corresponding to $\phi \to \phi^i\phi^j$ are obtained as

$$\mathcal{L} = -\frac{1}{2} e^K (K_{\phi}W_{ij} - 2\Gamma_{\phi}^{ik}W_{jk})^* g^{\phi\phi^*} W_{\phi\phi} \phi^i\phi^j + h.c.,$$  \hspace{1cm} (9)

where $g^{\phi\phi^*}e^{K/2}W_{\phi\phi} \equiv m_\phi$. Therefore, the decay into the right-handed sneutrino pair is the dominant decay channel for producing a pair of matter scalars, if it is allowed kinematically. The decay width is calculated as

$$\Gamma(\phi \to \tilde{\nu}_R\tilde{\nu}_R) = \frac{\lambda^2}{128\pi} N_f \sqrt{1 - \frac{4M^2}{m^2} \frac{M^2}{m^2} \frac{m^3_\phi}{M_{Pl}^2}}.$$  \hspace{1cm} (10)

The Polonyi field may decay into three-body final states such as $\phi \to \phi^i\chi^j\chi^k$ and $\phi \to \phi^i\phi^j\phi^k$. The former process occurs through the interactions as

$$\mathcal{L} = -\frac{1}{2} e^K (K_{\phi}W_{ijk} - 3\Gamma_{\phi}^{i\ell}W_{jk\ell})^* \phi^i\chi^j\chi^k + h.c.,$$  \hspace{1cm} (11)

and the three-scalar final state process is obtained by

$$\mathcal{L} = -\frac{1}{6} e^K (K_{\phi}W_{ijk} - 3\Gamma_{\phi}^{i\ell}W_{jk\ell})^* g^{\phi\phi^*} W_{\phi\phi} \phi^i\phi^j\phi^k + h.c.$$  \hspace{1cm} (12)
However, from the Kähler potential (4), the parentheses in eq. (11) and (12) vanish. The three-body decay process, then, does not occur.

In addition to these tree-level decay processes, the Polonyi field can also decay into the gauge supermultiplets through the anomaly-mediation effect [19], even if $\phi$ does not have any direct couplings to the gauge sector. Taking account of only the strong coupling, we obtain [19]

$$\Gamma_{\phi \rightarrow gg}^{\text{AM}} = \Gamma_{\phi \rightarrow \tilde{g} \tilde{g}}^{\text{AM}} \simeq \frac{9\alpha_s^2}{128\pi^3} \lambda^2 \frac{m_\phi^3}{M_{Pl}^2}. \quad (13)$$

If the decay processes $\phi \rightarrow \nu_R \nu_R$ and $\phi \rightarrow \tilde{\nu}_R \tilde{\nu}_R$ are not allowed kinematically, the dominant contribution to reheat the Universe comes from the anomaly-induced decay eq. (13). However, it is unlikely that such a decay width provides sufficiently high reheating temperature, since it is suppressed by a loop factor. On the other hand, if $\phi$ can decay into the right-handed (s)neutrino pair, there is a possibility that the process of decay into the right-handed (s)neutrino pair is more efficient than the anomaly-induced one. The possibility is that we will tune the right-handed (s)neutrino mass $M$ to maximize the decay width. In that case, the total decay width of $\phi$ comes from eq. (7) and eq. (10), that is,

$$\Gamma_{\text{tot}} \simeq \Gamma(\phi \rightarrow \nu_R \nu_R) + \Gamma(\phi \rightarrow \tilde{\nu}_R \tilde{\nu}_R)$$

$$= \frac{N_f}{32\pi} \sqrt{\frac{1 - 4M^2}{m_\phi^2}} \lambda \left( \frac{5}{4} - \frac{2M^2}{m_\phi^2} \right) M^2 \frac{m_\phi^3}{M_{Pl}^2}. \quad (14)$$

When the right-handed (s)neutrino mass $M$ satisfies the relation $M \simeq 0.38 m_\phi$, the total decay width is maximized as

$$\Gamma_{\text{tot}}^{\text{max}} \simeq 8.9 \times 10^{-3} N_f \alpha_s^2 \frac{m_\phi^3}{M_{Pl}^2}. \quad (15)$$

From eq. (15), the reheating temperature, $T_R(\phi)$, after $\phi$ decay is

$$T_R(\phi) \equiv \left( \frac{90}{\pi^2 g_*} \right)^{1/4} \sqrt{\Gamma_{\text{tot}}^{\text{max}} / M_{Pl}}$$

$$= 1.9 \times 10^{-3} \text{GeV} \lambda \sqrt{N_f} \left( \frac{g_*}{10} \right)^{-1/4} \left( \frac{m_\phi}{10^5 \text{GeV}} \right)^{3/2}, \quad (16)$$

where $g_*$ is the effective degrees of freedom of the radiation at the reheating. In Ref. [21], it was shown that in order to reproduce the observed abundance of $^4\text{He}$, the reheating temperature should be higher than about $4 - 7$ MeV [25] for the hadronic branching ratio
We find that $m_{\phi}$ should be heavier than $3 \times 10^5 \text{ GeV}$ for $\lambda = N_f = 1$ so that the reheating temperature survives this bound. Notice that in the absence of decay channels into right-handed (s)neutrino pair, the decay is dominated by eq. (13). In this case, the Polonyi mass should be heavier than about $2 \times 10^6 \text{ GeV}$. This implies that the gravitino mass is heavier than $10^6 \text{ GeV}$, which is out of the region of the low-energy supersymmetry.

Let us next consider a more stringent constraint imposed by the neutralino LSP relic abundance. Neutralino LSPs are produced by subsequent decay of the right-handed sneutrino, as well as through the anomaly-mediated decay eq. (13). Since the branching ratio of $\phi$ into neutralino LSPs is $\mathcal{O}(1)$, they are so abundant that the annihilation among them becomes effective. The annihilation process will cease when the Hubble parameter becomes comparable to the annihilation rate

$$\langle \sigma_{\text{ann}} v_{\text{rel}} \rangle n_{\text{LSP}} \simeq H(T_R),$$

where $\sigma_{\text{ann}}$ is the annihilation cross section of two LSPs, $v_{\text{rel}}$ their relative velocity, $\langle \cdots \rangle$ represents the thermal average, and $n_{\text{LSP}}$ is the number density of the LSPs. Thus, the yield of the LSPs at the $\phi$ decay can be estimated as

$$Y_{\text{LSP}} \simeq \frac{H(T_R)}{\langle \sigma_{\text{ann}} v_{\text{rel}} \rangle s} = \frac{1}{4} \left( \frac{90}{\pi^2 g_* (T_R)} \right)^{1/2} \frac{1}{\langle \sigma_{\text{ann}} v_{\text{rel}} \rangle T_R M_{\text{Pl}}}. \tag{18}$$

When we consider the case where the wino is the LSP, the annihilation process is most effective. The annihilation cross section is obtained as

$$\langle \sigma_{\text{ann}} v_{\text{rel}} \rangle = \frac{g_2^4}{16 \pi m_{\text{LSP}}^2} \frac{1}{(2 - x_W)^2}, \tag{19}$$

where $g_2$ is the $SU(2)$ gauge coupling constant, $m_{\text{LSP}}$ the LSP mass, and $x_W \equiv m_W^2/m_{\text{LSP}}^2$ with $W$ boson mass $m_W$. From eq. (18) and eq. (19), we can compute the ratio of the LSP mass density to the entropy density:

$$m_{\text{LSP}} n_{\text{LSP}} s \simeq 1.9 \times 10^{-9} \text{ GeV} \frac{1}{\lambda \sqrt{N_f}} \frac{(2 - x_W)^2}{(1 - x_W)^{3/2}} \left( \frac{m_{\text{LSP}}}{380 \text{ GeV}} \right)^3 \times \left( \frac{g_2}{10} \right)^{-1/4} \left( \frac{m_{\phi}}{3 \times 10^5 \text{ GeV}} \right)^{-3/2}, \tag{20}$$

where we have set the wino mass equal to 380 GeV. The reason is that since the Polonyi decay into the gravitino pair must be forbidden kinematically, gravitinos have to be heavier than $1.5 \times 10^5 \text{ GeV}$. Such a heavy gravitino induces the wino mass about $m_W \simeq 380 \text{ GeV}$ via
the anomaly mediated SUSY breaking effect \[23\]. It is convenient to write eq. (20) in terms of the Ω parameter which is defined by the ratio of the LSP mass density to the critical mass density,

$$\Omega_{\text{ann}} h^2 \simeq 0.53 \frac{1}{\lambda \sqrt{N_f}} \frac{(2 - x_W)^2}{(1 - x_W)^{3/2}} \left( \frac{m_{\text{LSP}}}{380 \text{ GeV}} \right)^3 \left( \frac{g_*}{10} \right)^{-1/4} \left( \frac{m_\phi}{3 \times 10^5 \text{ GeV}} \right)^{-3/2}, \tag{21}$$

where \( h \simeq 0.72 \) is the Hubble constant in units of 100 km/Mpc/s. The contours of eq. (21) in \( m_{\text{LSP}} - \lambda \) plane are shown in fig. 1. Recent WMAP observations \[24\] suggest that the density parameter of the dark matter be \( \Omega_{\text{DM}} h^2 = 0.105^{+0.007}_{-0.013} \) (68 % C.L.). Therefore, from fig. 1, in order to avoid the LSP overclosure, \( \lambda \gtrsim 20 \) even for the wino-like LSP. Since \( \lambda \) is the coupling constant between the Polonyi field and the matter fields, it is natural to expect that it is of order unity. Thus, even if the LSP is the wino, it cannot explain the present dark matter abundance. This conclusion also applies to the case where the LSP is bino- and higgsino-like one because the annihilation cross section is even smaller.

In summary, we have reconsidered the cosmological implications of the heavy Polonyi field with the mass \( \mathcal{O}(10) \text{ TeV} \). To avoid the heavy gaugino mass, the gauge kinetic function is assumed to be independent of the Polonyi field. In such a case, even when the Polonyi field can decay into the right-handed (s)neutrino pair and we have tuned the right-handed (s)neutrino mass \( M \) to maximize the total decay width, we found that the Polony mass well exceeds 100 TeV in order to obtain sufficiently high reheating temperature compatible with the standard BBN. In the absence of this decay channel, the Polony field dominantly decays into gauge multiplets through the anomaly mediation effect. In this case, however, the Polony mass has to be heavier than about 1000 TeV. We also discussed the neutralino LSP abundance produced by the Polonyi decay. We found that, if the neutralino LSP is stable, avoidance of the LSP overclosure requires \( \lambda \gtrsim 20 \) even when the LSP is a wino-like neutralino and its annihilation process is most efficient. This result implies that even for the wino-like LSPs, its abundance produced by the Polonyi decay cannot account for the present dark matter abundance. Other types of the neutralino LSPs would also be too abundant to be consistent with the WMAP observations. All in all, the explicit computation presented here makes the Polony problem even worse, and thus one should probably consider supersymmetry breaking scenarios in the absence of the Polonyi-type field.
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FIG. 1: Contours of the density parameter $\Omega_{\text{LSP}} h^2$ for the wino-like LSP drawn in $m_{\text{LSP}} - \lambda$ plane. Two lines represent $\Omega_{\text{LSP}} h^2 = 0.1, 0.3$, from the above.