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On the Solution to the Polonyi Problem with No-Scale Type Supergravity * †

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Abstract

We study the solution to the Polonyi problem in the framework of no-scale type supergravity. In such a model, Polonyi field can weigh as $O(10\text{TeV})$ and decay just before the big-bang nucleosynthesis. It is shown that in spite of a large entropy production by the decay of the Polonyi field, one can naturally explain the present value of the baryon-to-entropy ratio, $n_B/s \sim (10^{-10} - 10^{-11})$ if the Affleck-Dine mechanism for baryogenesis works. It is pointed out, however, that there is another cosmological problem related to the abundance of the lightest superparticles produced by the decay of the Polonyi field.

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1 Introduction

$N = 1$ supergravity [1] is not only regarded as an effective field theory of superstring below the Planck scale, but also provides a natural framework for the origin of the soft supersymmetry (SUSY)-breaking terms. Most of the supergravity models, however, contain a light massive boson ϕ (Polonyi field) with the mass m_ϕ of order the gravitino mass $m_{3/2}$ [2, 3, 4], which is responsible for the spontaneous SUSY breaking. The Polonyi field ϕ couples only gravitationally to the light particles and hence the lifetime of ϕ is very large as

$$\tau_\phi \simeq \Gamma_\phi^{-1} \sim \left(N \frac{m_\phi^3}{M_P^2} \right)^{-1}, \tag{1}$$

where Γ_ϕ is the decay rate of the Polonyi field, $M_P = \sqrt{8\pi} M \simeq 1.2 \times 10^{19}$GeV the Planck mass, and N the number of the decay modes. (In the following calculations, we take $N = 100$.) Then, the Polonyi field is expected to decay when the temperature of the universe becomes very low. The reheating temperature T_R due to the decay of the Polonyi field is given by

$$T_R \sim 1\text{MeV} \left(\frac{m_\phi}{10\text{TeV}} \right)^{3/2}. \tag{2}$$

This fact leads to a serious cosmological difficulty (so-called Polonyi problem) [2, 3, 4]. Under quite general assumptions, the Polonyi field ϕ takes an amplitude of order M at the end of inflation, and subsequently it starts oscillation and dominates the energy density of the universe until it decays. If the decay of the Polonyi field occurs after or during the big-bang nucleosynthesis, it most likely destroys one of the successful scenarios in the big-bang cosmology, that is the nucleosynthesis. Furthermore, the decay of the Polonyi field releases a tremendous amount of entropy and dilutes primordial baryon asymmetry much below what is observed today. Especially, the important point is that we cannot solve this problem even if we assume an inflation, which is the crucial difference between the Polonyi problem and another serious cosmological difficulty in $N = 1$ supergravity, *i.e.* the gravitino problem [5, 6, 7, 8, 9, 10].

It has been pointed out [11] that the first problem can be solved by raising the Polonyi mass m_ϕ (or equivalently the gravitino mass $m_{3/2}$) up to $O(10\text{TeV})$ so that the reheating temperature T_R by the decay of the Polonyi field becomes larger than $O(1\text{MeV})$. Then, the nucleosynthesis may re-start after the decay of the Polonyi field. This solution favors strongly "no-scale type" supergravity [12],[1] since the gravitino mass can be taken

[1]In the original no-scale supergravity model [13, 14], Polonyi field acquires a mass of the order of

$O(10\text{TeV})$ without diminishing the original motivation of SUSY as a solution to the hierarchy problem [16, 17]. Namely, we can raise the gravitino mass while keeping all masses of SUSY particles in observable sector to be $O(100\text{GeV})$.

Here, we stress that the second problem can be also solved if the Affleck-Dine mechanism [18] for baryogenesis works in the early universe [19]. However, we point out another cosmological problem that the lightest superparticles (LSPs) produced via the Polonyi decay are extremely abundant [19, 20]. As a result, their energy density, if stable, overcloses the universe unless the reheating temperature due to the Polonyi decay is sufficiently high. This fact gives us a lowerbound on the reheating temperature after the decay of the Polonyi field.

The organization is as follows. In the next section, we show how the baryon asymmetry of the universe can be explained if we assume the Affleck-Dine mechanism for baryogenesis. In section 3, we calculate the mass density of LSP due to the decay of the Polonyi field, and constrain the reheating temperature in the framework of the minimal SUSY SU(5) model. Section 4 is devoted to discussion.

2 Polonyi problem and the Affleck-Dine mechanism

The Affleck-Dine mechanism [18] for baryogenesis is based on the fact that there are some combinations of squark $\tilde{q}$ and slepton $\tilde{l}$ fields for which the scalar potential vanishes identically when SUSY is unbroken. After SUSY breaking, these flat-direction fields acquire masses m_χ of order 100GeV. One of these flat directions χ is assumed to have a large initial value χ_0 which is assumed to be about the grand unified theory (GUT) scale $M_{GUT} \sim 10^{16}$GeV or the gravitational scale. It has been shown [18] that the decay of the coherent oscillation mode of such a field χ can generate a large baryon-to-entropy ratio $\sim O(1)$ under the presence of tiny baryon-number violating operators such as $(m_S/M_{GUT})\tilde{q}\tilde{q}\tilde{q}\tilde{l}$ (with m_S being the scale of the SUSY breaking parameter in the observable sector, which is assumed to be $m_S \sim O(100\text{GeV})$).

We now compute how large baryon asymmetry can be obtained if we combine the Affleck-Dine mechanism with the Polonyi problem. For this purpose, it is convenient to use the fact that n_B/ρ_ϕ is independent of time since the baryon number is approximately conserved in the regime we consider. (Here, n_B is the baryon number density and ρ_ϕ the

$m_{3/2}^2/M$ [15] which is much smaller than the gravitino mass $m_{3/2}$. However, in the "no-scale type" supergravity model studied in Ref. [12], the mass of the Polonyi field is at the order of the gravitino mass.

mass density of the Polonyi field.) Then,

$$\frac{m_\chi n_B}{\rho_\phi} = \text{const.} \tag{3}$$

We evaluate this when the Affleck-Dine field χ starts its oscillation. At this time,

$$\frac{m_\chi n_B}{\rho_\phi} \simeq \frac{m_\chi n_B}{\rho_\chi}\frac{\rho_\chi}{\rho_\phi} \simeq \eta_{B0}\frac{\rho_\chi}{\rho_\phi} \simeq \eta_{B0}\left(\frac{\chi_0}{\sqrt{3}M}\right)^2. \tag{4}$$

where ρ_χ is the mass density of the Affleck-Dine field and $\eta_{B0} \equiv (n_B/n_\chi)_{H\simeq m_\chi}$ with n_χ being the number density of χ. In deriving Eq.(4), we have used $H \simeq \sqrt{\rho_\phi}/\sqrt{3}M$, and $\rho_\chi = m_\chi^2\chi_0^2$. On the other hand, we evaluate the same quantity given in Eq. (3) at the decay time of the Polonyi field ϕ

$$\frac{m_\chi n_B}{\rho_\phi} \simeq \frac{4}{3}\frac{m_\chi n_B(T_R)}{s(T_R)T_R}. \tag{5}$$

Equating Eq. (4) and Eq. (5), we get

$$\frac{n_B}{s} \simeq \frac{1}{4}\eta_{B0}\frac{T_R}{m_\chi}\left(\frac{\chi_0}{M}\right)^2 \sim 10^{-5}\eta_{B0}\left(\frac{T_R}{1\text{MeV}}\right)\left(\frac{100\text{GeV}}{m_\chi}\right)\left(\frac{\chi_0}{M}\right)^2. \tag{6}$$

With Eq.(6), one may explain the observed value $n_B/s \sim (10^{-10} - 10^{-11})$ taking $\chi_0 \sim M_{GUT}$, $T_R \sim 1\text{MeV}$, and $\eta_{B0} \sim O(1)$.[2]

In our case, the dilution factor D is given by

$$D \sim \frac{T_R}{m_\chi}\left(\frac{\chi_0}{M}\right)^2 \sim 10^{-5}\left(\frac{T_R}{1\text{MeV}}\right)\left(\frac{100\text{GeV}}{m_\chi}\right)\left(\frac{\chi_0}{M}\right)^2, \tag{7}$$

which is much larger than that derived in the previous work [11]. For example, the dilution factor given in Ref. [11] is $O(10^{-14})$ for the case $T_R \sim 1\text{MeV}$, which is about 10^{-9} times smaller than our result with $m_\chi \sim 100\text{GeV}$ and $\chi_0 \sim M$. This discrepancy originates to the fact that the amplitude of the Polonyi field has already decreased by a large amount at the decay time of the Affleck-Dine field. In Ref. [11], this effect is not taken into account, and hence the dilution factor given in Ref. [11] is underestimated.

[2]It has been pointed out that the Affleck-Dine mechanism for baryogenesis may result in too large baryon number fluctuation in the case of chaotic inflation [21]. However, such a difficulty can be solved if we adopt a larger value of the initial amplitude of the Affleck-Dine field; $\chi_0 \sim M$. In that case, we have to choose $\eta_{B0} \sim 10^{-5}$.

3 Mass density of LSP

Let us now turn to discuss a new cosmological difficulty in the present solution to the Polonyi problem. The decay of the Polonyi field produces a large number of superparticles, which promptly decay into LSPs. The number density of LSP produced by the decay, $n_{\rm LSP,i}$, is of the same order of that of the Polonyi field $n_\phi \equiv \rho_\phi/m_\phi$. Just after the decay of the Polonyi field, the yield variable for LSP, $Y_{\rm LSP}$, which is defined by the ratio of the number density of LSP to the entropy density s, is given by

$$\begin{aligned} m_{\rm LSP} Y_{\rm LSP} &\simeq \frac{\rho_\phi}{s} \simeq \frac{m_{\rm LSP}\rho_{\rm LSP,i}}{m_\phi s} \sim \frac{m_{\rm LSP} T_R}{m_\phi} \\ &\sim 10^{-5}{\rm GeV}\left(\frac{m_{\rm LSP}}{100{\rm GeV}}\right)\left(\frac{T_R}{1{\rm MeV}}\right)\left(\frac{10{\rm TeV}}{m_\phi}\right), \end{aligned} \tag{8}$$

where $\rho_{\rm LSP,i}$ is the mass density of LSP just after the decay of the Polonyi field. If LSP is stable and the pair annihilation of LSP is not effective, $Y_{\rm LSP}$ is conserved until today. On the other hand, the ratio of the critical density ρ_c to the present entropy density s_0 is given by

$$\frac{\rho_c}{s_0} \simeq 3.6 \times 10^{-9} h^2 \ {\rm GeV}, \tag{9}$$

where h is the Hubble constant in units of 100km/sec/Mpc. Comparing Eq.(8) with Eq.(9), we see that LSP overcloses the universe in the wide parameter region for $m_{\rm LSP}, m_\phi$ and T_R which we are concerned with.

If the pair annihilation of LSP takes place effectively, its abundance is reduced to

$$\frac{n_{\rm LSP}}{s} \simeq \left.\frac{H}{s\langle\sigma_{\rm ann} v_{\rm rel}\rangle}\right|_{T=T_R}, \tag{10}$$

where $\sigma_{\rm ann}$ is the annihilation cross section, $v_{\rm rel}$ is the relative velocity, and $\langle\cdots\rangle$ represents the average over the phase space distribution of LSP. Comparing Eq.(9) with Eq.(10), we obtain a lowerbound on the annihilation cross section,

$$\langle\sigma_{\rm ann} v_{\rm rel}\rangle \gtrsim 3 \times 10^{-8} h^{-2} {\rm GeV}^{-2} \left(\frac{m_{\rm LSP}}{100{\rm GeV}}\right)\left(\frac{100{\rm MeV}}{T_R}\right), \tag{11}$$

in order that the mass density of LSP does not overclose the universe.

As we can see, constraint (11) becomes severer as the reheating temperature T_R decreases, and hence we obtain a lowerbound on T_R. Here, we derive the constraint on T_R in the framework of minimal SUSY SU(5) model [23, 24], which is shown in Appendix A. We first solve RGEs based on the minimal SU(5) model with the no-scale boundary conditions, and determine the mass spectrum and mixing matrices of the superparticles. Notice

that we only investigate the parameter space which is not excluded by the experimental or theoretical constraints. The crucial constraints are as follows;

- Higgs bosons H_f and $\bar{H}_f$ have correct vacuum expectation values; $\langle H_f\rangle^2 + \langle \bar{H}_f\rangle^2 \simeq (174\text{GeV})^2$ and $\tan\beta = \langle H_f\rangle/\langle \bar{H}_f\rangle$.
- Perturbative picture is valid below the gravitational scale.
- LSP is neutral.
- Sfermions (especially, charged sleptons) have masses larger than the experimental lower limits [25].
- The branching ratio for Z-boson decaying into neutralinos is not too large [26].

One remarkable thing is that *LSP almost consists of bino which is the superpartner of the gauge field for* $U(1)_Y$ if we require that LSP is neutral. Therefore, in our model, the LSP mass m_{LSP} is essentially equivalent to the bino mass. Then, we calculate the annihilation cross section and determine the lowerbound on the reheating temperature from the following equation;

$$\left.\frac{H}{s\langle\sigma_{\text{ann}}v_{\text{rel}}\rangle}\right|_{T=T_R} \leq \frac{\rho_c}{s_0} \simeq 3.6h^2 \times 10^{-9}\text{GeV}. \tag{12}$$

Since LSP is most dominated by bino, it annihilates into fermion pairs. The annihilation cross section is given by [22]

$$\langle\sigma_{\text{ann}}v_{\text{rel}}\rangle = a + b\langle v^2\rangle, \tag{13}$$

where $\langle v^2\rangle$ is the average velocity of LSP, and

$$a \simeq \frac{32\pi\alpha_1^2}{27}\frac{m_t^2}{(m_{\tilde{t}_R}^2 + m_{\text{LSP}}^2 - m_t^2)^2}\left(1 - \frac{m_t^2}{m_{\text{LSP}}^2}\right)^{1/2}\theta(m_{\text{LSP}} - m_t), \tag{14}$$

$$b \simeq \frac{8\pi\alpha_1^2}{3}\sum_{m_f \leq m_{LSP}} Y_f^4\left\{\frac{m_{\text{LSP}}^2}{(m_{\text{LSP}}^2 + m_{\tilde{f}}^2)^2} - \frac{2m_{\text{LSP}}^4}{(m_{\text{LSP}}^2 + m_{\tilde{f}}^2)^3} + \frac{2m_{\text{LSP}}^6}{(m_{\text{LSP}}^2 + m_{\tilde{f}}^2)^4}\right\}. \tag{15}$$

Here, $\alpha_1^2 \equiv g_1^2/4\pi \simeq 0.01$ represents the fine structure constant for $\text{U}(1)_\text{Y}$, m_t the top-quark mass, Y_f the hypercharge of the fermion f, and $m_{\tilde{f}}$ the mass of the sfermion $\tilde{f}$. Notice that a- and b-terms correspond to s- and p-wave contributions, respectively. Taking $m_{\tilde{f}} \sim m_{\text{LSP}} \sim 100\text{GeV}$, the annihilation cross section given in Eq.(13) is at most

$3 \times 10^{-8}\mathrm{GeV}^{-2}$. Using this result in the inequality (11), we can see that the reheating temperature must be higher than about 100MeV even if $\langle v^2 \rangle \sim 1$. In fact, LSP is in kinetic equilibrium in the thermal bath [20], and hence its velocity is given by $O(T_R/m_{\mathrm{LSP}})$ which is much smaller than 1. Thus, we have severer constraint on T_R, as we will see below.

In Fig. 1, we show the lowerbound on the reheating temperature in the $\tan\beta$ vs. m_{LSP} plane. In the figures, large or small $\tan\beta$'s are not allowed since the Yukawa coupling constant for the top quark or bottom quark blows up below the gravitational scale for such $\tan\beta$'s. Furthermore, there also exists a lowerbound on the LSP mass. In the case where $\tan\beta \lesssim 20$, charged sfermions become lighter than the experimental limit if the LSP mass becomes lighter than $\sim$ 50GeV. On the other hand, for the large $\tan\beta$ case, unless the bino mass is sufficiently large, the lightest charged slepton becomes LSP. (Remember that the dominant component of LSP is bino.) Thus, the lowerbound on m_{LSP} is obtained. As we can see, the reheating temperature should be larger than about 100MeV, even for the case where $m_{\mathrm{LSP}} \sim 50\mathrm{GeV}$. The constraint becomes more stringent as m_{LSP} increases, since the masses of the superparticles which mediate the annihilation of LSP becomes larger as the LSP mass increases. If we translate the lowerbound on the reheating temperature into that of the Polonyi mass m_ϕ, we obtain $m_\phi \gtrsim 100\mathrm{TeV}$ (see Eq.(2)).

Finally, we comment on the accidental case where the annihilation process hits the Higgs pole in the s-channel. If the LSP mass is just half of the lightest Higgs boson mass, the LSP annihilation cross section is enhanced since LSP has small but nonvanishing fraction of higgsino component. If the parameters are well tuned, such a situation can be realized and the lowerbound of T_R decreases to $O(10\mathrm{MeV})$. However, we consider that such a scenario are very unnatural since a precise adjustment of the parameters is required in order to hit the Higgs pole. [3]

4 Discussion

Here, we proposed a solution to the Polonyi problem based on the no-scale type supergravity and the Affleck-Dine mechanism for baryogenesis. In our scenario, however, LSP may be overproduced due to the decay of the Polonyi field. From this fact, we obtained the lowerbound on the reheating temperature after the decay of the Polonyi field, which is given by $O(100\mathrm{MeV})$. As a result, the mass of the Polonyi field have to be larger than

[3] In the case where the annihilation process hits the pole of heavier Higgs bosons, the cross section is not enhanced so much, since the widths of the heavier Higgs bosons are quite large.

$O(100\text{TeV})$, which may raise a new fine-tuning problem [27, 28].

To cure this conflict in the case of $T_R \lesssim 10$ MeV, let us consider modifications of the minimal SUSY standard model (MSSM). One way is to extend the particle contents and provide a new, very light LSP. If the LSP is lighter than $O(10$ MeV$)$, we can see from Eq. (8) that the relic abundance does not exceed the critical density without invoking the annihilation. This is most easily realized in the minimal extension of the MSSM, where the superpartner of a singlet Higgs is contained in the neutralino sector. Another extension which has a light LSP is to incorporate the Peccei-Quinn symmetry. Then the superpartner of the axion, the axino, can be the LSP [29]. Indeed, it was shown in Ref. [30] that the axino becomes massless at the tree-level in the no-scale supergravity. Radiative corrections may give a small, model-dependent axino mass.[4] In the case of the axino mass $\sim$ 10MeV, the axino becomes a cold dark matter of the universe.

R-parity breaking is the other possibility to make our scenario cosmologically viable. In this case, the LSP is no longer stable, but decays to ordinary particles. If the lifetime τ_{LSP} of the LSP is shorter than 1 sec,[5] its decay does not upset the standard big-bang nucleosynthesis.

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A The model

In this appendix, we describe the model we use, *i.e.* the minimal SUSY SU(5) model [23, 24] with no-scale type boundary conditions. This model has three types of Higgs field; $H(\mathbf{5})$ and $\bar{H}(\mathbf{5}^*)$ which contain flavor Higgses H_f and $\bar{H}_f$, and $\Sigma(\mathbf{24})$ whose condensation breaks the SU(5) group into the gauge group of MSSM, $\mathrm{SU(3)_C} \times \mathrm{SU(2)_L} \times \mathrm{U(1)_Y}$. For the Higgs sector, the superpotential is given by

$$W = \frac{1}{3}\lambda \mathrm{tr}\Sigma^3 + \frac{1}{2}M_\Sigma \mathrm{tr}\Sigma^2 + \kappa \bar{H}\Sigma H + M_H \bar{H} H, \tag{16}$$

[4] A light axino can also be realized if one chooses a special form of superpotential [31].

[5] Such a small R-parity violation ($\tau_{LSP} \sim$ 1sec) is consistent with other phenomenological constraints [32].

where λ and κ are dimensionless constants, while M_Σ and M_H are mass parameters which are of the order of the grand unified theory (GUT) scale $M_{\rm GUT}(\sim 10^{16}{\rm GeV})$. Furthermore, the model also has the soft SUSY breaking terms;

$$\mathcal{L}_{\rm soft} = -\frac{1}{3}\lambda A_\Sigma {\rm tr}\Sigma^3 - \frac{1}{2}M_\Sigma B_\Sigma {\rm tr}\Sigma^2 - \kappa A_H \bar{H}\Sigma H - M_H B_H \bar{H}H + h.c., \tag{17}$$

where A_Σ, B_Σ, A_H and B_H are SUSY breaking parameters. Minimizing the Higgs potential, we find the following stationary point;

$$\langle\Sigma\rangle = \frac{1}{\lambda}\left\{M_\Sigma + 2\left(A_\Sigma - B_\Sigma\right) + O\left(\frac{A_\Sigma}{M_\Sigma}, \frac{B_\Sigma}{M_\Sigma}\right)\right\} \times {\rm diag}(2,2,2,-3,-3), \tag{18}$$

where the SU(5) is broken down to $\mathrm{SU(3)_C \times SU(2)_L \times U(1)_Y}$. Regarding this stationary point as the vacuum, we obtain MSSM as the effective theory below the GUT scale $M_{\rm GUT}$. Here, the masslessness of the flavor Higgses H_f and $\bar{H}_f$ is achieved by a fine tuning among several parameters; $M_H - 3\kappa M_\Sigma/\lambda \simeq \mu_H$, where μ_H is the SUSY-invariant Higgs mass in MSSM.

In the present model, the parameters in MSSM at the electroweak scale is obtained by solving renormalization group equations (RGEs). The boundary conditions on the parameters in the minimal SUSY SU(5) model are given at the gravitational scale M. Since we assume the no-scale type supergravity models, all the SUSY breaking parameters except for the gaugino mass vanish at the gravitational scale. From the gravitational scale to the GUT scale, the parameters follow the renormalization group flow derived from RGEs in the minimal SUSY SU(5) model. Then we determine the parameters in MSSM at the GUT scale through an appropriate matching condition between the parameters in the SUSY SU(5) model and those in MSSM. Finally, we use RGEs in MSSM from the GUT scale to the electroweak scale in order to obtain the low energy parameters.

As for the matching condition, we have a comment. In the stationary point (18), the mixing soft mass term of the two flavor Higgs bosons, $m_{12}^2 \bar{H}_f H_f$, is generated at the tree level, where m_{12}^2 is given by

$$m_{12}^2(M_{\rm GUT}) \simeq \left[\frac{6\kappa}{\lambda}(A_\Sigma - B_\Sigma)(A_H - B_\Sigma) - \mu_H B_H\right]_{\mu=M_{\rm GUT}}. \tag{19}$$

Since the mixing mass term depends on unknown parameters, λ and κ in Eq.(16), we regard m_{12}^2 as a free parameter taking account of the uncertainty of λ and κ in our analysis. Then, the low energy parameters are essentially determined by the gauge and Yukawa coupling constants and the following three parameters; the supersymmetric Higgs mass μ_H, the mixing mass of the two flavor Higgs bosons m_{12}^2, and the unified gaugino

mass.[6] However, it is more convenient to express these parameters by other physical ones. In fact, one combination of them is constrained so that the flavor Higgs bosons have correct vacuum expectation values; $\langle H_f \rangle^2 + \langle \bar{H}_f \rangle^2 \simeq (174\text{GeV})^2$. As the other two physical parameters, we use the mass of LSP, m_{LSP}, and the vacuum angle $\tan\beta \equiv \langle H_f \rangle / \langle \bar{H}_f \rangle$. Thus, once we fix m_{LSP} and $\tan\beta$, we can determine all the parameters in MSSM.[7]

References

[1] E. Cremmer, S. Ferrara, L. Grardello and A. van Proeyen, *Nucl. Phys.* **B212** (1983) 413.

[2] G.D. Coughlan, W. Fischler, E.W. Kolb, S. Raby and G.G. Ross, *Phys. Lett.* **B131** (1983) 59.

[3] T. Banks, D.B. Kaplan and A.E. Nelson, *Phys. Rev.* **D49** (1994) 779.

[4] B. de Carlos, J.A. Carlos, F. Quevedo and E. Roulet, *Phys. Lett.* **B318** (1993) 447.

[5] S. Weinberg, *Phys. Rev. Lett.* **48** (1982) 1303.

[6] J. Ellis, E. Kim and D.V. Nanopoulos, *Phys. Lett.* **B145** (1984) 181.

[7] R. Juszkiewicz, J. Silk and A. Stebbins, *Phys. Lett.* **B158** (1985) 463.

[8] M. Kawasaki and K. Sato, *Phys. Lett.* **B189** (1987) 23.

[9] T. Moroi, H. Murayama and M. Yamaguchi, *Phys. Lett.* **B303** (1993) 289.

[10] M. Kawasaki and T. Moroi, *Prog. Theor. Phys.* **93** (1995) 879.

[11] J. Ellis, D.V. Nanopoulos and M. Quiros, *Phys. Lett.* **B174** (1986) 176.

[12] H. Murayama, H. Suzuki, T. Yanagida and J. Yokoyama, *Phys. Rev.* **D50** (1994) 2356.

[13] J. Ellis, C. Kounnas and D.V. Nanopoulos, *Phys. Lett.* **B143** (1984) 410.

[6] In fact, parameters in MSSM slightly depend on the parameters in the SUSY GUT such as λ, κ and so on. In our numerical calculation, we ignore the effects of these parameters on the renormalization group flow.

[7] Yukawa coupling constants are determined so that the fermions have correct masses. The gauge coupling constants are also fixed so that their correct values at the electroweak scale are reproduced.

[14] J. Ellis, K. Enqvist and D.V. Nanopoulos, *Phys. Lett.* **B147** (1984) 99.

[15] J. Ellis, C. Kounnas and D.V. Nanopoulos, *Nucl. Phys.* **B241** (1984) 406.

[16] L. Maiani, *Gif-sur-Yvette Summer School on Particle Physics*, (Natl. Inst. Nucl. Phys. Part. Phys., Paris, 1979).

[17] M.J.G. Veltman, *Acta Phys. Pol.* **B12** (1981) 437.

[18] I. Affleck and M. Dine, *Nucl. Phys.* **B249** (1985) 361.

[19] T. Moroi, M. Yamaguchi and T. Yanagida, *Phys. Lett.* **B342** (1995) 105.

[20] M. Kawasaki, T. Moroi and T. Yanagida, preprint LBL-37715 (hep-ph/9509399).

[21] J. Yokoyama, *Astroparticle Phys.* **2** (1994) 291.

[22] K.A. Olive and M. Srednicki, *Phys. Lett.* **B230** (1989) 78.

[23] S. Dimopoulos and H. Georgi, *Nucl. Phys.* **B193** (1981) 150.

[24] N. Sakai, *Z. Phys.* **C11** (1981) 153.

[25] Particle Data Group, *Phys. Rev.* **D50** (1994) 1173.

[26] L3 Collaboration, *Phys. Lett.* **B350** (1995) 109.

[27] G.A. Diamandis, J. Ellis, A.B. Lahanas and D.V. Nanopoulos, *Phys. Lett.* **B173** (1986) 303.

[28] P. Binetruy and M.K. Gaillard, *Phys. Lett.* **B168** (1986) 347.

[29] K. Rajagopal, M.S. Turner and F. Wilczek, *Nucl. Phys.* **B358** (1991) 447.

[30] T. Goto and M. Yamaguchi, *Phys. Lett.* **B276** (1992) 103.

[31] E.J. Chun, J.E. Kim and H.P. Nilles, *Phys. Lett.* **B287** (1992) 123.

[32] B.A. Campbell, S. Davidson, J. Ellis and K.A. Olive, *Phys. Lett.* **B256** (1991) 457.

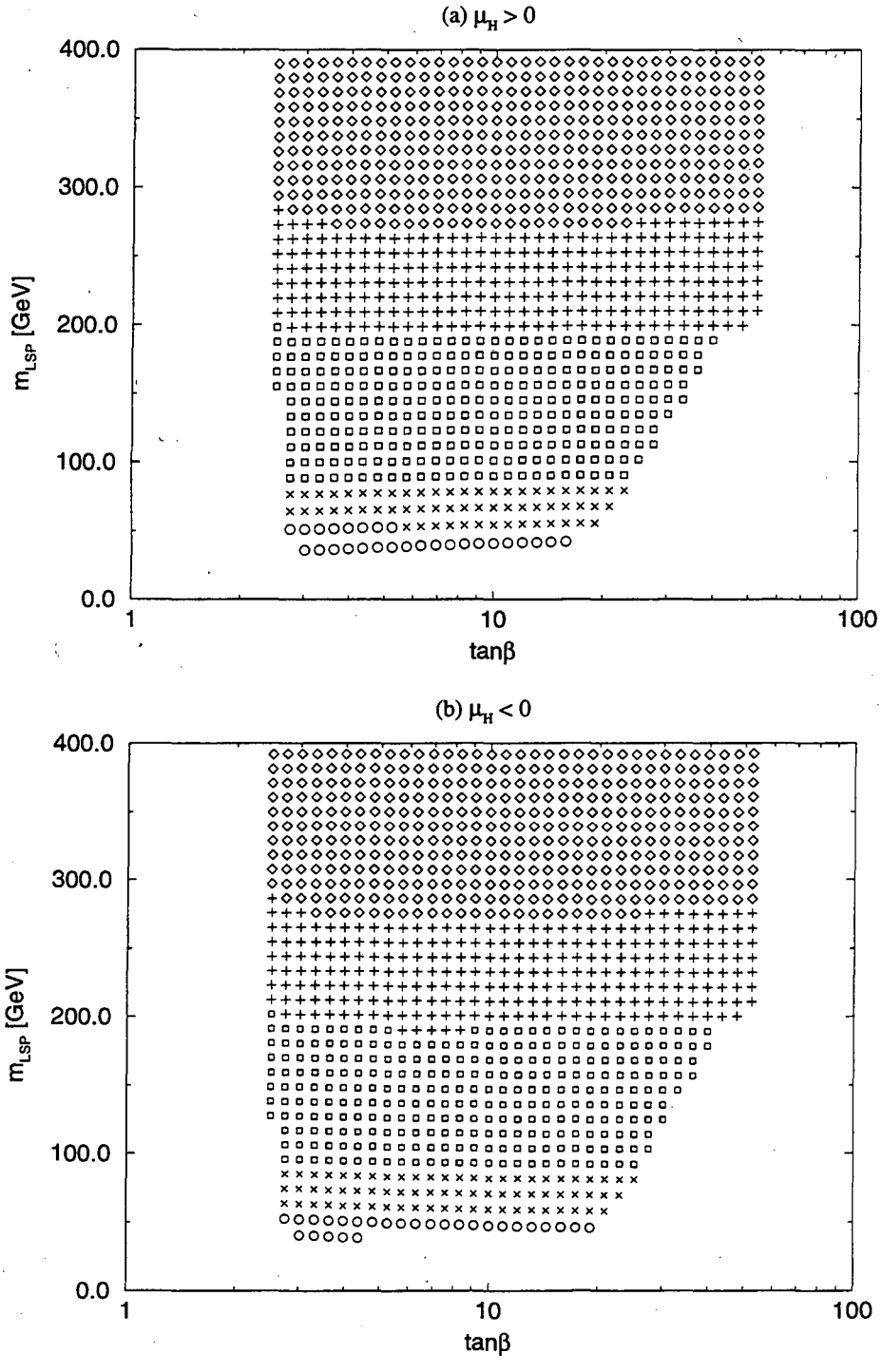

Figure 1: Lowerbound on T_R is shown in $\tan\beta$ vs. $m_{\rm LSP}$ plane. The meaning of each mark is as follows; $\circ$: 100MeV $\leq T_R \leq$ 500MeV, $\times$: 500MeV $\leq T_R \leq$ 1GeV, $\square$: 1GeV $\leq T_R \leq$ 5GeV, $+$: 5GeV $\leq T_R \leq$ 10GeV, $\diamond$: 10GeV $\leq T_R \leq$ 50GeV. The sign of the SUSY-invariant Higgs mass μ_H is taken to be (a) $\mu_H > 0$, and (b) $\mu_H < 0$.

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