Cosmological Light Element Bound on R Parity Violating Parameters

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Abstract

In R parity violating theories, there does not exist the stable lightest supersymmetric particle (LSP). The LSP, $\chi$, of MSSM decays to lighter ordinary particles with dimension 6 effective R violating interaction terms. Since the lifetime of $\chi$ can be sufficiently long for small R-violating couplings, it can affect the standard nucleosynthesis scenario. This constraint gives an upper limit for the lifetime of $\chi$ in the small $\tau_\chi$ region; $\sim \tau_\chi > 10^6$ s. This translates to the sum of squared R-violating couplings, for $m_\chi = 60$ GeV in the photino (bino) limit, $(0.12(0.05)|\lambda_{i,j,k}|^2 + 0.31(0.07)|\lambda'_{i,j\neq 3,k}|^2 + 0.04(0.04)|\lambda'_{i,3,k}|^2 + 0.23(0.12)|\lambda''_{i,j,k}(i < j, k \neq 3)|^2) > 7.7 \times 10^{-24}$. 
I. INTRODUCTION

Supersymmetric extension of the standard model is beautiful and probably needed toward a solution of the gauge hierarchy problem [1]. However, it introduces three outstanding problems: (i) the R-parity violating superpotential terms [2], (ii) the µ problem [3], and (iii) the flavor changing neutral currents [4]. There exist many ideas to remedy these problems.

In this paper, we concentrate on the R-parity violation and its probable effects in cosmology. The problem in the minimal supersymmetric standard model is that it allows the trilinear superpotential terms. For decay processes, let us consider the trilinear R-violating couplings defined by

$$W = \frac{1}{2} \lambda_{ijk} L^i L^j \bar{E}^k + \lambda'_{ijk} L^i Q^j \bar{D}^k + \frac{1}{2} \lambda''_{ijk} \bar{D}^i \bar{D}^j \bar{U}^k$$  \hspace{1cm} (1)

where the superfields are denoted as the upper case Roman letters. There can be the bilinear R-violating couplings, $\mu_i L_i H_2$, which can be rotated away at the superpotential level by redefining $H_1$ and $L_i$ fields. But if these bilinear parameters are given at the electroweak scale with soft terms present, one should worry about the correct minimum for the electroweak vacuum, implying that $\mu_i$ are physical parameters [5]. But this subtle point is not of relevance in our study of the neutralino decay. If both $\lambda'$ and $\lambda''$ terms are present, the stability of proton is a function of $\lambda' \lambda''$. To remove one of these terms for proton stability, some employ discrete symmetries [6]. But the most widely employed discrete symmetry is the R-parity, defined as $R = (-1)^{3B - L + 2S}$ where B is the baryon number, L is the lepton number and S is the spin of the particle. The fields in the standard model have $R = 1$ while their superpartners have $R = -1$. Namely superfields in Eq. (1) carries R charge –1, and all terms in Eq. (1) are forbidden by the R-parity conservation assumption. In addition, if R-parity is conserved, there exists an absolutely stable LSP (lowest mass supersymmetric particle) which carries $R = -1$ and cannot decay to ordinary particles due to the R conservation. This stable LSP can be a candidate for the constituent of the missing mass of the universe.

The reason one considers the terms given in Eq. (1) is that they are not forbidden by the gauge group and supersymmetry alone. Therefore, phenomenologically it is important to know bounds on the strengths of the coupling parameters $\lambda$, $\lambda'$ and $\lambda''$. The most stringent laboratory bounds on these R-violating couplings are summarized in Table 1 [7]. However, a more severe bound comes from the lower limit of the lifetime of proton as mentioned above. If we take the universal strength for the couplings, then $\lambda$’s must be less than $10^{-12}$ from the proton stability condition. With this strong constraint the search for new effects in accelerator experiments may be futile.

Table 1. The best phenomenological limits on R violating couplings for $M_{\tilde{f}} = 100$ GeV [7].

| $\lambda_{ijk}$ | $\lambda'_{ijk}$ | $\lambda''_{ijk}$ |
|-----------------|-----------------|-----------------|
| (121) $< 0.05$ | (111) $< 0.001$ | (112) $< 10^{-6}$ |
| (122) $< 0.05$ | (112) $< 0.02$  | (113) $< 10^{-5}$ |

From the Table, we note that there exist several regions of the couplings allowed by the laboratory experiments, for $\lambda$, $\lambda'$ and $\lambda''$. Generally, particle phenomenology gives an upper bound on the couplings in view of nonobservation of rare processes.
On the other hand, cosmological constraints give different kinds of bounds on the R-violating couplings. Some of the allowed regions are reviewed in Ref. [8]. In this paper, we study the deuterium dissociation after nucleosynthesis, and present a lower bound on the sum of squares of couplings. There can be a window for the allowed strength of the couplings. We try to estimate the lower bound of the couplings from deuterium dissociation as best as we can, including every possible factor in the formulae. In fact, this topic was touched upon earlier [9], and there are a few papers regarding which process (D or HE4 dissociation) in the gravitino decay gives a better bound [10]. In this paper, the lower bound on the R-violating couplings from the study of neutralino decay is expressed as an upper bound on the lifetime of the LSP (≡ χ). Of course, if the lifetime is sufficiently long then our analysis should not give any deuterium destruction since deuteriums are too much separated. This will happen sometime after photon recombination. However, the study shows that the upper limit of the lifetime falls in $10^6$ s region, and we are safe from the recombination phenomena. But if the lifetime is somewhat longer than the age of universe then the LSP has not decayed yet, then the only constraint is its mass density not exceeding the energy density of the universe. This gives another region of coupling constants, given as an upper bound, which is not considered here.

In Sec. 2, the decay rate of the neutralino is calculated. In Sec. 3, we calculate the decoupling temperature and obtain $T_D \sim O(1)$ GeV. In Sec. 4, these results are used to estimate the energy density of the neutralino and deuterium dissociation effect from photons arising in the neutralino decay. This study gives a bound on the sum of squares of the R-violating couplings. Sec. 5 is a conclusion.

II. DECAY RATE

With the superpotential given in Eq. (1), the LSP of MSSM $\chi$, decays to ordinary particles. For example, with the above R-parity violating terms, one can write, for instance, a dimension 6 effective Lagrangian

$$-\frac{\lambda_{ijk}}{M_f^2} (gN_2 + \frac{1}{3}g_1 N_1) \int d^4 \theta U^{a ij} \chi \int d^2 \theta E^a D^k + \text{h.c.}$$

where $a$ is the color index, and $N_1$ and $N_2$ are defined by

$$\bar{\tilde{\chi}} = N_1 \tilde{B} + N_2 \tilde{W} + N_3 \tilde{H}_1^0 + N_4 \tilde{H}_2^0. \quad (2)$$

For photino, one may write

$$\chi = \tilde{\gamma} + F^{\mu \nu} \sigma_{\mu \nu} \theta + \cdots.$$ 

But we use the vector multiplet notation in the Wess-Zumino gauge

$$-\theta \sigma^{\mu} \bar{\theta} A_\mu(x) + i \theta \bar{\theta} \tilde{\gamma}(x) - i \bar{\theta} \theta \tilde{\gamma}(x) + \frac{1}{2} \theta \bar{\theta} \bar{\theta} D(x). \quad (3)$$

For the decay $\chi \rightarrow \bar{u} + d + e^+$, there exist 9 diagrams among which one superfield diagram is shown in Fig. 1. In total, there exist 1,080 (1,350) decay modes for the decay of $\chi$ to any final states, excluding (including) the final state $t\bar{t}$ pair.
Assuming that $\chi$ is much heavier than b-quark, but lighter than t-quark, and a universal sfermion mass, we obtain the decay width of $\chi$ as:

$$
\Gamma_{\text{tot}} = \frac{g^2}{256\pi} \frac{m_\chi^5}{M_f^4} \left[ \left( \sum |\lambda_{ijk}|^2 \right) \left\{ \frac{1}{8} N_2^2 + \frac{3}{8} N_1^2 \tan^2 \theta_w \right\} 
+ \left( \sum_{j \neq 3} |\lambda'_{ijk}|^2 \right) \left\{ \frac{7}{12} N_2^2 + \frac{5}{12} N_1^2 \tan^2 \theta_w \right\} 
+ \left( \sum_{ik} |\lambda''_{ijk}|^2 \right) \left\{ \frac{1}{8} N_2^2 + \frac{7}{24} N_1^2 \tan^2 \theta_w - \frac{1}{2} N_1 N_2 \tan \theta_w \right\} 
+ \left( \sum_{k \neq 3} |\lambda''_{ijk}|^2 \right) \left\{ N_2^2 \tan^2 \theta_w \right\} \right].
$$

For $M_f = 1$ TeV and $m_\chi = 30$ GeV, the numerical factor in front of the square bracket is $1.8 \times 10^{15}$ s$^{-1}$. If the largest value of the R-violating coupling is of order $3 \times 10^{-8}$, then the LSP lifetime falls in the seconds region, which may affect the standard scenario of nucleosynthesis.

### III. DECOUPLING TEMPERATURE

The effect of a late decaying unstable (light) neutrino has been studied two decades ago [12]. Several years later, the cosmological effect of a late decaying heavy particle has attracted a great deal of attention, due to the possible existence of 100 GeV scale gravitino [13]. The gravitino lifetime can be very long due to the Planck mass suppression of the coupling constant. The lifetime of the LSP in R-violating theories can be very long due to the smallness of the R-violating parameters and three body final states. But their roles in cosmology, being a late decaying particle, are similar.

The decoupling temperature of the gravitino is close to the Planck scale, so the inflation might have diluted them almost completely. But after reheating, the gravitinos can be produced. Even though the number density of these thermally produced gravitinos are negligible compared to the number density of photon, its cosmological effect is dramatic since the gravitino decay injects huge energy in the cosmic soup [13]. On the other hand, the decoupling temperature $T_D$ of LSP is at a few GeV scale as shown below. If the reheating temperature after inflation is greater than the LSP decoupling temperature $T_D$, then the LSP number density is enormous compared to the thermally produced gravitino number density. This consideration will give a lower bound on the R-violating coupling strengths. Of course, there exists another region of allowed R-violating couplings, the upper bound of couplings, so that the LSP can be practically considered as a stable particle in cosmology.

The decoupling temperature of $\chi$, $T_D$, is determined by a competition to the contribution to the number density of $\chi$ from the expansion rate of the universe, i.e. the depletion of $\chi$ through $H \cdot (\text{the } \chi \text{ number density})$, and from the $\chi$ destruction rate, $\sigma \cdot (\text{relative velocity}) \cdot (\text{useful number densities})^2$. The destruction rate is a function of sfermion mass for $\chi\chi \rightarrow f \bar{f}$. However, if all sfermion masses are much greater than the $Z$ boson mass and if $\chi$ has a

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1For a formal expression of matrix element squared, see [11].

2This number is a practical lower bound for a neutralino mass in MSSM.
significant Higgsino component, then the $Z$ exchange diagram can dominate. Since most superpartner masses are unknown at present, we illustrate for the latter possibility that the LSP has significant Higgsino component, i.e. $N_{3,4}$ of Eq. (3) are relatively large $|N_{3,4}/\sqrt{N_1^2 + N_2^2}| > M_Z^2/M_f^2$, the production of $\chi$ is dominated by the intermediate $Z$ boson diagram, and hence the destruction cross section for $T_D$ estimate is given by

$$\sigma(\chi\chi \rightarrow e^+e^-) = \frac{G_F^2 s \beta}{24\pi^2} (1 - 4x_w + 8x_w^2) \left|\frac{M_Z^2}{s - M_Z^2 + iM_Z\Gamma_Z}\right|^2 (N_3^2 - N_4^2)^2$$ \hspace{1cm} (5)$$

where $s$ is the total energy squared $x_w = \sin^2 \theta_w \sim 0.232$ and $\beta$ is the velocity of $\chi$, $\beta = \nu/c = \sqrt{1 - 4m_\chi^2/s}$. We are interested in the mass region $m_\chi < 100$ GeV. For $2m_b < m_\chi < M_Z$, the channels $\chi\chi \rightarrow e^+e^-, \mu^+\mu^-, \tau^+\tau^-, \nu_e \bar{\nu_e}, \nu_\mu \bar{\nu_\mu}, \nu_\tau \bar{\nu_\tau}, u\bar{d}, d\bar{d}, s\bar{s}, c\bar{c}$, and $b\bar{b}$ are open. In this region, we must consider all these final states, and Eq. (5) is changed as

$$(1 - 4x_w + 8x_w^2) \rightarrow (21 - 40x_w + \frac{160}{3}x_w^2) \approx 14.6.$$  

For a given temperature, the Boltzmann distribution of a massive particle is folded to the above cross section. In the literature [14], the thermal average regarding the relative velocity of neutralino has been properly taken into account. Thus, the Hubble expansion rate and the reaction rate are related to find the decoupling temperature, $<\sigma v_{12}> n_\chi \sim H$ where $v_{12} = |v_1 - v_2|$ is the relative velocity,

$$\frac{G_F^2 m_\chi T}{2\pi^2} (21 - 40x_w + \frac{160}{3}x_w^2) \left\{ \frac{M_Z^4}{(4m_\chi^2 - M_Z^2)^2 + M_Z^2 \Gamma_Z^2} \cdot \left( N_3^2 - N_4^2 \right)^2 \cdot g \left(\frac{m_\chi T}{2\pi}\right)^{3/2} e^{-\frac{m_\chi}{T}} \right\} = N_F^{1/2} \frac{T^2}{M_p} \sqrt{\frac{8\pi^3}{45}}.$$ \hspace{1cm} (6)$$

We can take $g = 2$ and $N_F = 494/8$ in the region of our interest. Then Eq. (6) gives a decoupling temperature through

$$\frac{m_\chi}{T_D} \approx \ln \left( \frac{G_F^2 m_\chi^{5/2} T_D^{1/2} M_p g(21 - 40x_w + \frac{160}{3}x_w^2)M_Z^2(N_3^2 - N_4^2)^2}{\sqrt{N_F^{8\pi^3/45}(2\pi)^{5/2}(4m_\chi^2 - M_Z^2)^2 + M_Z^2 \Gamma_Z^2}} \right)$$ \hspace{1cm} (7)$$

For $m_\chi = 30$ GeV and $|N_3^2 - N_4^2| \sim 0.1$, we obtain $T_D \sim 1.3$ GeV. In Table 2–4, we list the decoupling temperature factor $d$, which is defined as

$$d = \frac{m_\chi}{T_D},$$ \hspace{1cm} (8)$$

as a function of the neutralino mass, not only for the Higgsino dominated cases [14] and also for the gaugino dominated cases [12]. In these Tables, $M_f = 300, 600$ and $1,000$ GeV are considered, and we distinguish four cases: [i] $N_1 = 0.99, N_2 = 0.14, N_3 = N_4 = 0$, [ii] $N_1 = 0.9, N_2 = 0.44, N_3 = N_4 = 0$, [iii] $N_1 = 0.5, N_2 = 0.2, N_3 = 0.84, N_4 = 0$, [iv] $N_1 = 0.2, N_2 = 0.05, N_3 = 0, N_4 = 0.97$. As one can see, the Higgsino dominated cases

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3For a more concrete discussion, see, Ref. [15]
give lower decoupling temperatures than the gaugino dominated cases. This is because for most parameter space the Higgsino dominated cases give larger cross sections than gaugino dominated cases. This implies that the Higgsino dominated cases give smaller relic densities than the gaugino dominated cases. For the $\chi$ dark matter possibility, $\chi$ may be dominated by gaugino.

If $|N_{3,4}| \ll 1$ so that $\chi$ is mostly gaugino, then the $t$ and $u$ channel sfermion exchange diagrams are more important ones for establishing equilibrium. This happens when $N_{1,2} < M_Z^2/M_f^2 \sim 0.1$ if $M_f > 300$ GeV. The mass dependence of the cross section times the relative velocity is roughly $0.03 m_\chi^2/M_f^2$ and the decoupling temperature is $\sim 1.5$ GeV for the bino dominated case. For the detailed result, see Table 2. We also note that the neutralino decoupling temperature is insensitive to $\tan \beta$. In this calculation, we used the expression given in Drees and Nojiri [13], which employed the partial wave analysis. In this estimate, we set all the sfermion masses equal, for simplicity. We also assumed

$$\frac{m_{A,h_1,h_2} \Gamma_{A,h_1,h_2}}{m_\chi^2} \simeq \frac{1}{1,500}, \quad \frac{m_{A,h_1,h_2}}{m_\chi} \simeq \frac{1}{3}$$

(9)

where $m_{A,h_1,h_2}$ are the masses of the pseudoscalar, the lighter scalar and the heavier scalar, respectively, in the Higgs sector, $\Gamma_{A,h_1,h_2}$ are their decay rates. But it turns out that the cross sections and the decoupling temperatures are generally insensitive to these parameters. In the calculation we also take the limit $m_f/m_\chi = 0$ where $m_f$ is the final fermion mass. We also disregard the mixing of left and right sfermions because these are not well known.

The gaugino and Higgsino components can decay to a fermion and an anti-fermion through the $s$-channel Higgs boson exchange of the three kinds ($A, h_1, h_2$). The Higgs boson coupling to the fermion pair is proportional to $m_d/M_W \cos \beta$ and $m_u/M_W \sin \beta$ for the down-type and up-type quarks, respectively. So the Higgs boson coupling is important only for the $t$ and $b$ (if $\tan \beta$ is very large) quarks. Since we are interested in the light neutralino case, we do not consider the following decay modes: $\chi \chi \to t\bar{t}, W^+W^-, ZZ, Zh_{1,2}, ZA, W^\pm H^\mp, h_{1,2}h_{1,2}, AA, h_{1,2}A$, and $H^+H^-$ where $H^\pm$ is the charged Higgs particle.

If one consider the sfermion dominated region, $T_D$ falls in the $O(1$ GeV$)$ region for the interesting range of $M_f$ and $m_\chi$. Then the number density of neutralinos below the decoupling temperature is about $\sim 10^{-9}$ times the photon number density, which is enormous. If it is stable, neutralinos can form the missing mass of the universe.

IV. LIFETIME BOUNDS

If $\chi$ decays after 1 s, it causes problems during the nucleosynthesis era. From Eq. (4), the lifetime $\tau_\chi$ is about $10^{-15}/|\lambda|^2$ s. In the radiation dominated era, the cosmic time is given by the cosmic temperature as

$$\frac{m_\chi}{112}$$

[4] In view of the gauge hierarchy solution through supersymmetry, superpartner masses are not far from the electroweak scale. Under this circumstance, the LSP mass might be below 100 GeV.
\[ t = \left( \frac{90}{32\pi^3N} \right)^{1/2} \frac{M_{Pl}}{T^2}. \]  
\hspace{1cm} (10)

Long after the nucleosynthesis era, one constraint that the neutralino energy density does not exceed the energy of relativistic particles is

\[ m_\chi n_\chi(\tau_\chi) < \left( \frac{\pi^2}{30} \right) N_{\tau_\chi} T_{\tau_\chi}^4 \]  
\hspace{1cm} (11)

where \( T_{\tau_\chi} \) is the temperature of the universe when neutralinos decay and \( N_{\tau_\chi} \) is the effective degrees at the time of \( \chi \) decay. Using \( N_{\tau_\chi} \sim 2.5 \), since the expected temperature at the decay time is less than \( \sim \) keV, we obtain if \( \tau_\chi \) is greater than 1 s,

\[ \frac{m_\chi}{T_{\tau_\chi}} < 4.0 \cdot 10^8 \]  
\hspace{1cm} (12)

where we used an approximate relation \( T_D \simeq m_\chi/20 \). Using Eq. (6), we estimate

\[ \tau_{sec} < 4.8 \cdot 10^{10} \left( \frac{\text{GeV}}{m_\chi} \right)^2. \]  
\hspace{1cm} (13)

where \( \tau_{sec} \) is the lifetime of \( \chi \) in units of seconds. From Eq. (4), we estimate \( \lambda \)'s for the bino (photino) case if the couplings are universal,

\[ |\lambda| > 3.7(6.4) \times 10^{-15} \left( \frac{m_\chi}{\text{GeV}} \right). \]  
\hspace{1cm} (14)

However, this is not the most stringent bound for \( \tau_\chi > 1 \) s. When \( \chi \) decays long after the nucleosynthesis era, the universe containing light nuclei such as D, \(^3\)He, \(^4\)He and \(^7\)Li might be affected. Since the observed abundances of these light nuclei are in good agreement with the conventional calculation of these elements through nucleosynthesis, we must ensure that the light ordinary particles produced by \( \chi \) decay should not dissociate these precisely formed light nuclei.

The most stringent bound comes from the study of the D dissociation \([16,17]\). The calculation of Ref. \([16]\) can be expressed as

\[ (m_\chi f) \frac{\beta n_\chi}{E_* n_e} \leq 1 \]

where \( f = 0.7 \), \( E_* = 100 \ \text{MeV} \), \( n_\chi/n_\gamma \sim 3 \cdot 10^{-9} \), and \( \beta = 0.23 \) at \( T \simeq 100 \ \text{eV} \) at the time of \( \chi \) decay. If the above bound is not satisfied, \( \chi \) should decay so that its number density changes drastically. This implies that the parameters do not satisfy the above bound, then the lifetime of \( \chi \) must be less than 1 s, i.e. it decays before D synthesis. Assuming that \( n_e \sim n_B \) (the baryon density), we obtain for a given \( \delta_B = n_B/n_\gamma \)

\[ (m_\chi f)_{\text{GeV}} < \{130 \ (\text{for } \delta_B \sim 10^{-10}), \ 13 \ (\text{for } \delta_B \sim 10^{-9}), \ 1.3 \ (\text{for } \delta_B \sim 10^{-8})\}. \]

\(^5\)Depending on the allowed R-violating couplings, \( \beta \) can be different. Here, we take an eyeball number. Anyway, a better estimate is given below.
Namely, for interesting regions of the baryon asymmetry and the $\chi$ mass the condition on the lifetime of $\chi$ requires a detailed study.

However, this simple calculation \[16\] does not include the effect of $e^+e^-$ production from scattering on the background radiation by the high energy photon debrises from $\chi$ decay. This process has been added for a better estimation of $n_e$ and the resulting D dissociation rate \[17\]. The Compton scattering and the $e^+e^-$ production from scattering on the background radiation is comparable at the critical temperature $T_*$ with the critical photon energy $E_*$, satisfying

$$E_*T_* \simeq \frac{1}{50} \text{(MeV)}^2.$$  \hspace{1cm} (15)

At the cosmic temperature $T_*$, there are roughly $10^{-9}$ fraction of photons having energy $w \geq 25T_*$ which can be used to produce $e^+e^-$ pairs. Since these energetic photon number is comparable to the electron number, the probability to scatter off the electron through Compton scattering is comparable to the probability to produce $e^+e^-$ pairs from background radiation, and gives the relation (14).

When the photon energy exceeds $E_*$, the probability of D dissociation is unimportant compared to the total cross section, since the probability is proportional to $\sigma_D/\sigma_{\text{tot}}$ where $\sigma_D$ is the D dissociation cross section and $\sigma_{\text{tot}}$ is the total cross section. But the total cross section decreases rapidly for the photon energy below $E_*$. In this case, the $e^+e^-$ produced by the scattering on the background radiation scatter off the low energy photons to transfer the energy to photons. Thus, \textit{there results some high energy photons with energy less than $E_*$ but high enough ($> 2.225$ MeV) to dissociate D.} These photons are the major source of destructing D. In Ref. \[17\] this effect has been taken into account. It is summarized for the maximum allowable lifetime of $\chi$, satisfying

$$2x_0\frac{\epsilon}{\epsilon_0} \int_{t_1}^{t_2} \Sigma(\epsilon_0, T)e^{-t/\tau_{\text{max}}} \frac{dt}{\tau_{\text{max}}} = 1$$ \hspace{1cm} (16)

where $x_0$ is $n_\chi/n_e$ at the time of $\chi$ decay, and $\epsilon$ is the energy of the electron. We can take $t_1$ as the cosmic time when the maximum secondary photon energy equals $E_*$ and $t_2$ as the cosmic time when the maximum scattered photon energy equals the threshold energy for D dissociation. The above formula assumes a scaling behavior of $\Sigma$

$$\Sigma(\epsilon, T) = \frac{\epsilon}{\epsilon_0} \Sigma(\epsilon_0, T).$$ \hspace{1cm} (17)

The scaling behavior is not valid for low energy photons in which case we use the original $\Sigma$,

$$\Sigma(\epsilon, T) = \int_{Q_D}^{w_1} \frac{\sigma_D}{\sigma_{KN} + \sigma_{pp}} 2 \left(1 - \frac{w}{w_m}\right) \frac{dw}{w_m}$$ \hspace{1cm} (18)

where $Q_D = 2.225$ MeV. For low energy electrons the maximum scattered photon energy is given by

$$w_m \simeq \frac{12\gamma^2 T}{1 + \frac{12\gamma T}{m_e}} \sim 12\gamma^2 T.$$ \hspace{1cm} (19)
In Eq. (18), $\sigma_{KN}$ is the Klein-Nishina formula [18] and $\sigma_{pp}$ is the pair production cross section off the proton. For an accurate calculation, we will use Eq. (17).

But for illustration, let us take the scaling law, Eq. (16), and calculate $\Sigma$ at $E_* \sim 100$ MeV, in which case $\epsilon_0 \sim 50$ MeV, and $T_* \sim 0.2$ keV. Then,

$$\Sigma(\epsilon_0, T) \sim \frac{(w_1 - Q_D)(2w_m - w_1 - Q_D)}{w_m^2} \left( \frac{\sigma_D}{\sigma_{KN} + \sigma_{pp}} \right)_{E_\gamma = w_m, E_e = \epsilon_0}.$$

(20)

In this case, $w_m \sim 24$ MeV, and Eq.(16) gives

$$\tau_{\text{max}} = \frac{2.24 \times 10^7 \text{ s}}{4.92 + \ln M_{\text{GeV}} - \ln \eta_9}$$

(21)

where $M_{\text{GeV}} = m_\chi/\text{GeV}$ and $\delta_B = 10^{-9} \eta_9$. For $m_\chi = 30$ GeV, the maximum allowable lifetime of $\chi$ is $2.7 \times 10^6$ s.

But in contrast to this simple illustration, the condition (16) is rather complicated, and we must solve it numerically. In Figures 2–9, we present the numerical solution of Eq. (16) for the maximum lifetime bounds of the neutralino as a function of the neutralino mass. We remind again that these figures are drawn with the assumption of the same sfermion mass. But note that the almost Higgsino case does not depend on the sfermion mass if it is much larger than the $Z$ boson mass. The gaugino dominated case depends on the Higgsino mass, and here our assumption of the universal sfermion mass is critical. Figures 2–4 correspond to $N_1 = 0.99, N_2 = 0.14, N_3 = N_4 = 0$, and Figures 5–7 correspond to $N_1 = 0.9, N_2 = 0.44, N_3 = N_4 = 0$. Fig. 8 and Fig. 9 correspond to $N_1 = 0.5, N_2 = 0.2, N_3 = 0.84, N_4 = 0$, and $N_1 = 0.2, N_2 = 0.05, N_3 = 0, N_4 = 0.97$, respectively. Note that the $N_i$ dependence is meager, i.e. only logarithmic.

For the gaugino dominated cases (Figs. 2–7), we translate this bound to the lower bound on the sum of squared couplings with appropriate coefficients. This condition gives a bound on the sum of squares of R-violating couplings for a 60 GeV photino-like (bino-like) neutralino and for $\delta_B \approx 5 \times 10^{-10}$,

$$\left( \sum_{0.12(0.05)} |\lambda_{i,j,k}|^2 + \sum_{0.31(0.07)} |\lambda'_{i,j\neq3,k}|^2 + \sum_{0.04(0.04)} |\lambda'_{i,3,k}|^2 \right) + \sum_{0.23(0.12)} |\lambda''_{i,j,k}(i < j, k \neq 3)|^2 > 7.7 \cdot 10^{-24}.$$  

(22)

Thus a generic bound on the R-violating couplings, if the couplings are of the same order, is about $2.8 \times 10^{-12}$, which is a better bound than Eq. (14).

V. CONCLUSION

We have studied the R-parity violation effects on the dissociation of light nuclei, in particular to deuterium. We tried to include every factor in the formulae so that the estimate is reliable within a factor of a few. From this study, we obtain that the lifetime of the neutralino LSP must be shorter than $\sim 2 \times 10^8$ s, which translates to the sum of squared R-violating couplings being greater than $10^{-22}$. Therefore, the lifetime of the neutralino between $2 \times 10^6$ s and $\sim$ a few $\times 10^{17}$ s is forbidden.
If the strengths of R-violating couplings are comparable, which is possible in some spontaneous breaking scenario of R-violation, the lower bound in the large coupling region is $\sim 10^{-12}$. Such a small coupling implies a very small ratio for the mass parameters, $\tilde{v}/M_P \sim 10^{-12}$ where $\tilde{v}$ is the scale of spontaneous R-parity violation. Thus the lower limit of R-parity violating scale is far below the intermediate scale for the gaugino condensation and the axion decay constant. But this region is almost forbidden by the proton stability condition, which requires $\lambda' \lambda'' < 10^{-24}$. With the universal strength assumption on the R-violating couplings and to keep the $\lambda > 10^{-12}$ to be viable for some region of the coupling space, $\lambda'$ or $\lambda''$ must be required to vanish, e.g. from an additional symmetry.

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Table Caption

Table 2. The decoupling temperature factor $d$ for $M_{	ilde{f}} = 300$ GeV. The decoupling temperature is given by $T_D = m_\chi/d$. $N_i$ are defined in $\chi = N_1 \tilde{B} + N_2 \tilde{W} + N_3 \tilde{h}_1 + N_4 \tilde{h}_2$. The four cases for $N_i$ are [i] $N_1 = 0.99, N_2 = 0.14, N_3 = N_4 = 0$, [ii] $N_1 = 0.9, N_2 = 0.44, N_3 = N_4 = 0$, [iii] $N_1 = 0.5, N_2 = 0.2, N_3 = 0.84, N_4 = 0$, and [iv] $N_1 = 0.2, N_2 = 0.05, N_3 = 0, N_4 = 0.97$.

Table 3. The decoupling temperature factor $d$. Same as Table 2 except for $M_{\tilde{f}} = 600$ GeV.

Table 4. The decoupling temperature factor $d$. Same as Table 2 except for $m_{\tilde{f}} = 1,000$ GeV.
Table 2

| $m_{\chi}$ [GeV] | $\tan \beta$ | [i] | [ii] | [iii] | [iv] |
|------------------|--------------|-----|-----|------|-----|
| 30               | 2            | 20.25 | 20.32 | 32.83 | 33.40 |
| 30               | 10           | 20.25 | 20.32 | 32.83 | 33.40 |
| 30               | 30           | 20.25 | 20.32 | 32.84 | 33.40 |
| 30               | 50           | 20.25 | 20.32 | 32.85 | 33.40 |
| 60               | 2            | 21.49 | 21.56 | 29.26 | 29.83 |
| 60               | 10           | 21.49 | 21.56 | 29.27 | 29.83 |
| 60               | 30           | 21.49 | 21.56 | 29.36 | 29.83 |
| 60               | 50           | 21.49 | 21.56 | 29.52 | 29.83 |
| 100              | 2            | 22.24 | 22.31 | 29.49 | 30.06 |
| 100              | 10           | 22.24 | 22.31 | 29.50 | 30.06 |
| 100              | 30           | 22.24 | 22.31 | 29.54 | 30.06 |
| 100              | 50           | 22.24 | 22.31 | 29.64 | 30.06 |
| 150              | 2            | 22.62 | 22.69 | 28.97 | 29.53 |
| 150              | 10           | 22.62 | 22.69 | 28.97 | 29.53 |
| 150              | 30           | 22.62 | 22.69 | 29.03 | 29.53 |
| 150              | 50           | 22.62 | 22.69 | 29.17 | 29.53 |
Table 3.

| $m_\chi$ [GeV] | $\tan \beta$ | Column i | Column ii | Column iii | Column iv |
|----------------|--------------|----------|-----------|------------|-----------|
| 30             | 2            | 17.58    | 17.64     | 32.83      | 33.40     |
| 30             | 10           | 17.58    | 17.64     | 32.83      | 33.40     |
| 30             | 30           | 17.58    | 17.64     | 32.84      | 33.40     |
| 30             | 50           | 17.58    | 17.64     | 32.85      | 33.40     |
| 60             | 2            | 18.90    | 18.97     | 29.26      | 29.83     |
| 60             | 10           | 18.90    | 18.97     | 29.27      | 29.83     |
| 60             | 30           | 18.90    | 18.97     | 29.34      | 29.83     |
| 60             | 50           | 18.90    | 18.97     | 29.47      | 29.83     |
| 100            | 2            | 19.83    | 19.90     | 29.49      | 30.06     |
| 100            | 10           | 19.83    | 19.90     | 29.50      | 30.06     |
| 100            | 30           | 19.83    | 19.90     | 29.52      | 30.06     |
| 100            | 50           | 19.83    | 19.90     | 29.57      | 30.06     |
| 150            | 2            | 20.49    | 20.56     | 28.97      | 29.53     |
| 150            | 10           | 20.49    | 20.56     | 28.97      | 29.53     |
| 150            | 30           | 20.49    | 20.56     | 29.00      | 29.53     |
| 150            | 50           | 20.49    | 20.56     | 29.06      | 29.53     |
Table 4.

| $m_\chi$ [GeV] | $\tan \beta$ | i    | ii   | iii  | iv   |
|----------------|--------------|------|------|------|------|
| 30             | 2            | 15.60| 15.67| 32.83| 33.40|
| 30             | 10           | 15.60| 15.67| 32.83| 33.40|
| 30             | 30           | 15.60| 15.67| 32.84| 33.40|
| 30             | 50           | 15.60| 15.67| 32.86| 33.40|
| 60             | 2            | 16.94| 17.00| 29.26| 29.83|
| 60             | 10           | 16.94| 17.00| 29.27| 29.83|
| 60             | 30           | 16.94| 17.00| 29.34| 29.83|
| 60             | 50           | 16.94| 17.00| 29.47| 29.83|
| 100            | 2            | 17.90| 17.97| 29.49| 30.06|
| 100            | 10           | 17.90| 17.97| 29.50| 30.06|
| 100            | 30           | 17.90| 17.97| 29.52| 30.06|
| 100            | 50           | 17.90| 17.97| 29.56| 30.06|
| 150            | 2            | 18.65| 18.72| 28.97| 29.53|
| 150            | 10           | 18.65| 18.72| 28.97| 29.53|
| 150            | 30           | 18.65| 18.72| 28.97| 29.53|
| 150            | 50           | 18.65| 18.72| 29.03| 29.53|
FIG. 1. A superfield diagram for the gaugino component of $\chi$. The intermediate state is $Q$. There exists another diagram with the external $Q$ and $L$ lines exchanged.
FIG. 2. The maximum allowable lifetime of $\chi$ for $M_{\tilde{f}} = 300$ GeV.

$N_1 = 0.99, N_2 = 0.14, N_3 = N_4 = 0$.

FIG. 3. Same as Fig. 2 except for $M_{\tilde{f}} = 600$ GeV.
FIG. 4. Same as Fig. 2 except for $M_\tilde{f} = 1,000$ GeV.

FIG. 5. Same as Fig. 2 except for $N_1 = 0.9, N_2 = 0.44, N_3 = N_4 = 0$ and $M_\tilde{f} = 300$ GeV.
FIG. 6. Same as Fig. 5 except for $M_f = 600$ GeV.

FIG. 7. Same as Fig. 5 except for $M_f = 1,000$ GeV.
FIG. 8. The upper limit of $\tau_\chi$ for $N_1 = 0.5$, $N_2 = 0.2$, $N_3 = 0.84$, $N_4 = 0$. The $M_f$ dependence is negligible.

FIG. 9. Same as Fig. 8 except for $N_1 = 0.2$, $N_2 = 0.05$, $N_3 = 0$, $N_4 = 0.97$. 