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Baryogenesis with Superheavy Squarks

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Abstract

We consider a setup where R-parity is violated in the framework of split supersymmetry. The out-of-equilibrium decays of heavy squarks successfully lead to the generation of a baryon asymmetry. We restrict the R-parity violating couplings to the baryon number violating subset to keep the neutralino sufficiently stable to provide the dark matter. The observed baryon asymmetry can be generated for squark masses larger than 10^{11} GeV, while neutralino dark matter induces a stronger bound of 10^{13} GeV. Some mass splitting between left- and right-handed squarks may be needed to satisfy also constraints from gluino cosmology.

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1 Introduction

The general supersymmetric extension of the standard model (SM) contains baryon and lepton number violating interactions [1, 2]. Often a discrete symmetry, called R-parity, is imposed to eliminate these operators. In conventional supersymmetric models there are stringent bounds on many of the R-parity violating couplings induced by proton decay, neutrino masses, flavor violation, etc. The breaking of R-parity also has important cosmological implications. The baryon and lepton number violating interactions may erase a baryon asymmetry created in the early universe. They also induce the decay of the lightest supersymmetric particle (LSP). Only in the case of extremely small R-parity violating couplings does the LSP remain a viable dark matter candidate.

Most of these constraints are considerably relaxed in the context of split supersymmetry, where all scalars, except for a single (and finely tuned) Higgs boson, have masses much larger than the electroweak scale, $\tilde{m} \lesssim 10^{13}$ GeV [3–5]. For instance, bounds from proton decay have been discussed in ref. [6]. Taking $\tilde{m} = 10^{13}$ GeV, the product of baryon and lepton number violating couplings can be as large as 10^{-5}. In conventional supersymmetry it is bounded to be below 10^{-25} to guarantee a sufficiently long-lived proton.

The lepton number violating superpotential contains bilinear terms, $\mu'_i L_i H_2$. They lead to a tree-level mixing between the neutrinos and neutralinos which is not suppressed by the heavy mass scale of split supersymmetry. This mixing generates a neutrino mass at the tree-level which constrains the μ'-terms to be smaller than about 1 MeV [7]. It also induces the decay of the lightest neutralino, e.g. into a neutrino and a (virtual) Z boson. Keeping the neutralino lifetime larger than the age of the universe requires $\mu'_i \lesssim 10^{-8}$ eV. In contrast, neutralino decay via R-parity violating Yukawa couplings (trilinear terms) is indeed suppressed by the large sfermion masses [8]. The lepton number violating Yukawa couplings induce μ'-terms in the low energy effective action by quantum corrections even if they are set to zero at the tree-level [9]. Thus, an enormous amount of tuning is needed in this setting to keep the lightest neutralino sufficiently stable to provide the dark matter.

In this article we show how the baryon asymmetry of the universe could be generated by the out-of-equilibrium decays of heavy sfermions in the framework of split supersymmetry[2]. In this context, neutralino dark matter is an important link to tie the masses of the gauginos and higgsinos to the electroweak scale [3–5]. To avoid rapid neutralino decay by lepton number violating bilinears, we restrict R-parity violation to the baryon number violating operators.

There have been attempts to use R-parity violating interactions to generate the baryon asymmetry in weak scale supersymmetry [11]. However, these typically require some additional ingredients to provide the necessary departure from thermal equilibrium, such as non-thermal production of squarks at the end of inflation, late decay of gravitinos or axinos, etc. We will show that in split supersymmetry ther-

[2]See ref. [10] for an alternative proposal.

mal baryogenesis is successful for squark masses larger than about 10^{11} GeV. As it turns out, the neutralino lifetime induces a somewhat stronger constraint on the squark masses of about 10^{13} GeV. Gluino cosmology may require some mass splitting between left- and right-handed squarks.

The organization of the paper is as follows. In sec. 2 we discuss how a CP-asymmetry arises in the squark decays. It requires the interference of tree-level and two-loop diagrams. In sec. 3 we write down the Boltzmann equations, which describe the baryon production process. Their numerical solution together with some analytical approximations will be presented in sec. 4. Constraints induced by the gluino and neutralino lifetimes will be discussed in sec. 5. In sec. 6 we present our conclusions.

2 CP-asymmetry in squark decays

According to the discussion in the introduction we consider the following R-parity violating interactions in the superpotential [1, 2][3]

$$W_{\not R_P} = \lambda_{ijk} U_i^c D_j^c D_k^c, \tag{1}$$

where $\lambda_{ijk} = -\lambda_{ikj}$. These operators violate baryon number but not lepton number. The proton is therefore stable. The form of the superpotential (1) could, for instance, be guaranteed by lepton parity.

Since A-terms, gaugino masses and the μ-parameter are small compared to the split supersymmetry scale $\tilde{m}$ [3–5], scalar masses for U_i^c and D_i^c are the only soft terms relevant to our discussion. We can diagonalize them by supersymmetric rotations. Then flavor transitions can only be induced by the R-parity violating couplings (1) or by ordinary Yukawa couplings. We make the (conservative) assumption that the effects of the latter are small, and ignore them.

We assume a (somewhat) hierarchical spectrum of the SU(2) singlet ("right-handed") squarks[4]. Then baryogenesis is dominated by the decay of the lightest of these states, which we take to be the right-handed up squark $\tilde{u}$. It will turn out that this choice simplifies the Boltzmann equations, which describe the baryogenesis process. Taking another right-handed squark would change the final baryon asymmetry only by a factor of order 1.

The spectrum of the SU(2) doublet ("left-handed") squarks is not directly related to the baryogenesis process. If some of them are lighter than $\tilde{u}$, they only enter the Boltzmann equations as additional degrees of freedom, which carry part of the baryon number. Later on this possibility will turn out to be helpful to reconcile the constraints induced by the gluino and neutralino lifetimes. A somewhat

[3]Note that in the literature these couplings are usually denoted by λ''_{ijk}.

[4]For the SU(2) singlet states we use the notation $\tilde{u} \equiv (\tilde{u}^c)^*$ etc., so that they are counted with positive baryon number.

lighter left-handed squark can speed up the gluino decay through supergauge interactions, making it more easy to satisfy constraints from gluino cosmology [12]. The neutralino decay by R-parity violating couplings is still governed by the (heavier) right-handed squarks. Thus the neutralino can remain sufficiently long-lived to provide the dark matter of the universe.

At least four of the nine couplings of the superpotential (1) have to be present to allow for CP violation in the up squark decays. In the following we assume two of them to be λ_{112} and λ_{123}. Two other couplings have to be related to a different, heavier squark, which we take to be $\tilde{t}$. The first two operators induce baryon number violating decays of the up squark. At tree-level the partial decay widths are

$$\begin{aligned}\Gamma(\tilde{u} \to \bar{d}_1 \bar{d}_2) \equiv B_{12}\Gamma_D &= \frac{|\lambda_{112}|^2}{8\pi} m_{\tilde{u}} \\ \Gamma(\tilde{u} \to \bar{d}_2 \bar{d}_3) \equiv B_{23}\Gamma_D &= \frac{|\lambda_{123}|^2}{8\pi} m_{\tilde{u}},\end{aligned} \tag{2}$$

where Γ_D denotes the total decay width and $m_{\tilde{u}}$ the up squark mass. Note that the d quarks are also right-handed (SU(2) singlet) states. For simplicity we ignore a possible decay into the $d_1 d_3$ channel. The up squark also decays by the supergauge interaction at a rate

$$\Gamma(\tilde{u} \to u\tilde{g}) \equiv B_g \Gamma_D = \frac{2}{3}\alpha_s m_{\tilde{u}}. \tag{3}$$

Finally, the up squark can decay through its Yukawa coupling y_u. This contribution is suppressed by the small value of y_u and ignored in the following.

At the loop-level, the decay widths for $\tilde{u}$ and and its antiparticle $\tilde{u}^*$ become different in the presence of CP-violation

$$\begin{aligned}\Gamma(\tilde{u} \to \bar{d}_1 \bar{d}_2) &= \Gamma_D \left(B_{12} - \frac{\epsilon_{12}}{2}\right) \\ \Gamma(\tilde{u} \to \bar{d}_2 \bar{d}_3) &= \Gamma_D \left(B_{23} - \frac{\epsilon_{23}}{2}\right) \\ \Gamma(\tilde{u} \to \bar{u}\tilde{g}) &= \Gamma_D \left(B_g + \frac{\epsilon_g}{2}\right).\end{aligned} \tag{4}$$

Going to antiparticles, the sign of the CP-violating contributions are reversed. Since $\tilde{u}$ and its anti-particle $\tilde{u}^*$ have the same width,

$$\epsilon_g = \epsilon_{12} + \epsilon_{23} \tag{5}$$

holds.

A CP-asymmetry in the decay width is generated by the interference of tree-level and 2-loop amplitudes[5]. The diagram of fig. 1a leads to

$$\epsilon_{12} \propto \mathrm{Im}(\lambda_{112}\lambda^*_{123}\lambda^*_{312}\lambda_{323}) \propto \sin(\delta), \tag{6}$$

[5]Note that 1-loop contributions are suppressed by the small gaugino masses and A-parameters.

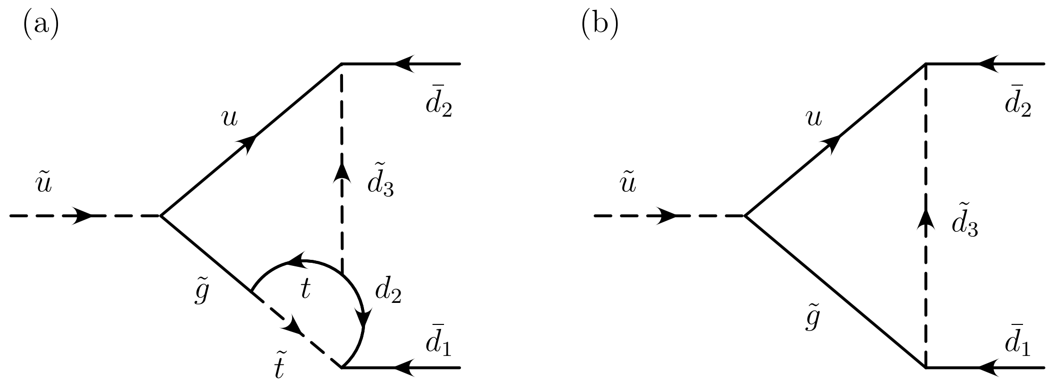

Figure 1: 2-loop contribution (a) to ϵ_{12} and its approximation by a 1-loop diagram (b).

where λ_{312} and λ_{323} are related to the (heavier) top squark, and δ defines the effective CP-phase. Instead of computing the 2-loop diagram, we approximate the effect of the stop by a non-diagonal gluino vertex

$$\sqrt{2}\theta_{13} g_3 \tilde{g} \tilde{d}_1^* d_3. \tag{7}$$

Computing the remaining 1-loop diagram of fig. 1b , we obtain for the CP-asymmetry

$$\epsilon_{12} = \frac{2}{3}\alpha_s \sin(\delta) B_{12}\theta_{13} \left|\frac{\lambda_{123}}{\lambda_{112}}\right| f\left(\frac{m_{\tilde{u}}^2}{m_{\tilde{d}_3}^2}\right). \tag{8}$$

The loop function is given by

$$f(x) = 2\left(1 - \frac{1}{x}\ln(1+x)\right) \approx x + \mathcal{O}(x^2). \tag{9}$$

Of course, there are more diagrams that contribute[6], so this only a rough estimate. Note that ϵ_{12} can be enhanced by taking the ratio $\lambda_{123}/\lambda_{112}$ to be large. The CP-asymmetry ϵ_{23} in the $d_2 d_3$ channel is obtained by $\delta \to -\delta$ and $m_{\tilde{d}_3} \to m_{\tilde{d}_1}$.

3 Boltzmann equations

The evolution of the baryon asymmetry is governed by a set of Boltzmann equations. We assume that kinetic equilibrium is maintained and approximate the phase-space densities by Maxwell-Boltzmann distributions,

$$f_i(E) = e^{-(E-\mu_i)/T}, \tag{10}$$

[6]The off-diagonal gluino vertex can also be generated by ordinary Yukawa couplings, with left-handed squarks running in the loop.

where μ_i is the chemical potential of the ith particle species. It describes the deviation from chemical equilibrium.

3.1 Up squark decay

The number density of right-handed up squarks is governed by the Boltzmann equation

$$\frac{dn_{\tilde{u}}}{dt} + 3Hn_{\tilde{u}} = -D_g - D_{12} - D_{23}. \tag{11}$$

Here $H = 1.66\sqrt{g_*}T^2/M_P$ denotes the Hubble parameter, with $M_P = 1.22 \times 10^{19}$ GeV. If all scalars are heavier than $\tilde{u}$, the number of degrees of freedom in the plasma is $g_* = 131.25$. For every left-handed squark doublet, which is light, g_* increases by 12. To first order in the CP-asymmetries, the collision terms for the different decay channels are given by

$$\begin{aligned}
D_g &= \gamma_D \left[n_{\tilde{u}} \left(B_g + \frac{\epsilon_g}{2} \right) - n_{\tilde{u}}^{\text{eq}} (1 + \xi_u) \left(B_g - \frac{\epsilon_g}{2} \right) \right] \\
D_{12} &= \gamma_D \left[n_{\tilde{u}} \left(B_{12} - \frac{\epsilon_{12}}{2} \right) - n_{\tilde{u}}^{\text{eq}} (1 - \xi_{d_1} - \xi_{d_2}) \left(B_{12} + \frac{\epsilon_{12}}{2} \right) \right] \\
D_{23} &= \gamma_D \left[n_{\tilde{u}} \left(B_{23} - \frac{\epsilon_{23}}{2} \right) - n_{\tilde{u}}^{\text{eq}} (1 - \xi_{d_2} - \xi_{d_3}) \left(B_{23} + \frac{\epsilon_{23}}{2} \right) \right],
\end{aligned} \tag{12}$$

where $\xi_i = \mu_i/T$. To obtain this result we approximated $\exp(\xi_i) \approx 1 + \xi_i$ and used $\xi_{\bar{i}} = -\xi_i$ for the antiquark chemical potentials. The quark chemical potentials are first order in the CP-asymmetries, so that to first order we can drop terms like $\xi\epsilon$. We used energy conservation to eliminate the quark distribution functions and CPT invariance to obtain the rates for the inverse decay processes. The equilibrium number density $n_{\tilde{u}}^{\text{eq}}$ is computed from the distribution function with vanishing chemical potential and includes three degrees of freedom. The thermal decay rate averaged with equilibrium distribution function is [13]

$$\gamma_D(z) = \frac{K_1(z)}{K_2(z)} \Gamma_D, \tag{13}$$

where K_1 and K_2 are Bessel functions and

$$z = \frac{m_{\tilde{u}}}{T}. \tag{14}$$

The evolution equation for $\tilde{u}^*$ is obtained by reversing the signs of the CP-violating contributions and replacing the quark chemical potentials by that of antiquarks, and vice versa.

It is useful to write the number densities of $\tilde{u}$ and $\tilde{u}^*$ as [13]

$$N_\pm = n_{\tilde{u}} \pm n_{\tilde{u}^*}. \tag{15}$$

Up to linear order in the CP-violating asymmetries, we obtain

$$\frac{dN_+}{dt} + 3HN_+ = -\gamma_D(N_+ - N_+^{\rm eq}) \tag{16}$$

The equilibrium density $N_+^{\rm eq} = 2n_{\tilde{u}}^{\rm eq} = (3T^3/\pi^2)z^2K_2(z)$ contains six degrees of freedom.

We normalize the number densities to the entropy density $s = (2\pi^2 g_* T^3/45)$, where g_* is again the number of degrees of freedom in the plasma, such that

$$Y_i = \frac{N_i}{s}. \tag{17}$$

Normalization by s removes the term proportional to the Hubble parameter from the Boltzmann equation. It is convenient to transform to the dimensionless variable z, using $dY/dt = (z/H(T = m_{\tilde{u}}))(dY/dz)$, to arrive at

$$\frac{dY_+}{dz} = -\beta K(Y_+ - Y_+^{\rm eq}), \tag{18}$$

where

$$\begin{aligned} \beta(z) &= z\frac{K_1(z)}{K_2(z)} \\ K &= \frac{\Gamma_D}{H(T = m_{\tilde{u}})}. \end{aligned} \tag{19}$$

Assuming that the decay rate is dominated by the gluino channel, we have $K \approx 1.5 \times 10^3$ $(10^{13}\ \mathrm{GeV}/m_{\tilde{u}})$, where the running strong coupling is $\alpha_s(10^{13}\ \mathrm{GeV}) = 0.035$. A large value of K signals that the up squark decay is close to equilibrium. Along the same lines we arrive at

$$\frac{dY_-}{dz} = -\beta K(Y_- - Y_+^{\rm eq}[B_g\xi_u - B_{12}(\xi_{d_1} + \xi_{d_2}) - B_{23}(\xi_{d_2} + \xi_{d_3})]). \tag{20}$$

We see that the inverse decays lead to a back reaction of the quark densities on the evolution of Y_-.

3.2 Evolution equations for the quark densities

Before computing the evolution of the quark densities we have to check what other than the R-parity violating interactions are relevant at temperatures around 10^{13}GeV which we are going to consider.

i) Strong sphalerons transform left-handed quarks into right-handed ones. Using the results of ref. [14] and taking into account the running of α_s, we find these processes to be in equilibrium for $T \lesssim 5 \times 10^{13}$ GeV (see also ref. [15]).

ii) Weak sphalerons [16] become fast at temperatures $T \lesssim 2 \times 10^{12}$ GeV.

iii) The scattering rate of Yukawa interactions has been estimated as $\Gamma_y \approx 0.02y^2T$ [17]. The top Yukawa coupling therefore is in equilibrium for $T \lesssim 10^{16}$ GeV. At $T \lesssim 1 \times 10^{12}$ GeV also the bottom and tau Yukawa interactions become relevant.

In the following we will assume the strong sphalerons and the top Yukawa interaction to be in equilibrium. All other Yukawa couplings and the weak sphalerons will be neglected. This approximation is valid in the temperature range 1×10^{12} GeV $\lesssim T \lesssim 5\times 10^{13}$ GeV and leaves the leptons entirely out of the baryon generation process.

The relevant operators are

$$\begin{aligned} &\tilde{g}\tilde{u}^*u,\ \tilde{u}d_1d_2,\ \tilde{u}d_2d_3,\ \tilde{g}\tilde{d}_3^*d_1, \\ &\bar{t}q_3H, \\ &uctd_1d_2d_3\bar{q}_1\bar{q}_1\bar{q}_2\bar{q}_2\bar{q}_3\bar{q}_3. \end{aligned} \tag{21}$$

The quark species c, q_1 and q_2 are only produced by strong sphalerons. Therefore the number densities are related as

$$2N_c = -N_{q_1} = -N_{q_2} \equiv 2N_q. \tag{22}$$

Here we use the number densities of particles minus antiparticles, i.e. $N_c = n_c - n_{\bar{c}}$etc. The operators (21) conserve four independent U(1) charges, which imply the constraints

$$\begin{aligned} N_t + N_{q_3} + N_q &= 0 \\ N_{d_1} - N_{d_2} + N_{d_3} - N_q &= 0 \\ N_- + N_u - N_{d_2} &= 0 \\ N_t + N_H - N_q &= 0. \end{aligned} \tag{23}$$

Taking the top Yukawa and strong sphaleron interactions to be in equilibrium forces the corresponding linear combinations of chemical potentials to vanish. What these relations mean in terms of number densities depends on the number of degrees of freedom (d.o.f.) related to each chemical potential. To be specific, we assume that one Higgs doublet is much heavier than $m_{\tilde{u}}$ and the supergauge interactions maintain the equilibrium between the Higgs and higgsino densities. The Higgs chemical potential then represents 12 d.o.f., like q_3. As already mentioned, we can allow for some left-handed squarks to be light, which we take to be $\tilde{q}_1$. Again supergauge interactions will equilibrate the q_1 and $\tilde{q}_1$ densities and effectively double the d.o.f. of q_1. The equilibrium conditions then take the form

$$\begin{aligned} -2N_t + N_{q_3} + N_H &= 0 \\ N_u + N_t + N_{d_1} + N_{d_2} + N_{d_3} - N_{q_3} + \left(3 + \frac{2}{k}\right) N_q &= 0, \end{aligned} \tag{24}$$

where $k = 1$ (2) if $\tilde{q}_1$ is heavy (light). Using eqs. (22), (23) and (24) we can express all quantities in terms of N_-, N_u and N_{d_1} as

$$\begin{aligned} N_{d_2} &= N_- + N_u \\ N_{d_3} &= \frac{3k+2}{5k+2}N_- + \frac{2k+2}{5k+2}N_u - N_{d_1} \\ N_q &= -N_{q_3} = N_H = -\frac{2k}{5k+2}N_- - \frac{3k}{5k+2}N_u \\ N_t &= 0. \end{aligned} \tag{25}$$

We can use these relations to express baryon number as

$$\begin{aligned} N_B &\equiv \frac{1}{3}(N_- + N_u + N_c + N_t + N_{d_1} + N_{d_2} + N_{d_3} + N_{q_1} + N_{q_2} + N_{q_3}) \\ &= \frac{7k+2}{5k+2}N_- + \frac{8k+2}{5k+2}N_u. \end{aligned} \tag{26}$$

Note that N_{d_1} does not enter here directly. This is an accident of the approximations we are using. For instance, it would no longer hold if the bottom Yukawa coupling were taken into account.

To compute the generated baryon asymmetry we need two evolution equations in addition to eqs. (18), (20). We take them to be

$$\begin{aligned} \frac{d}{dz}(Y_u - Y_c) &= D_g - S_{12} - S_{23} - T_{12} - T_{23} - S'_{12} - S'_{23} - T'_{12} - T'_{23} - \text{CP conj.} \\ \frac{d}{dz}(Y_{d_1} - Y_c) &= -D_{12} - S_{12} - T_{12} - S_{13} - S'_{12} - T_{13} - T'_{12} - \text{CP conj.} \end{aligned} \tag{27}$$

Subtracting Y_c in these equations removes the strong sphaleron rate from the right-hand side, which we have taken to be in equilibrium. The inverse decays D_{12}, D_{23} and D_g are given by eq. (12). We can use $n_u - n_{\bar{u}} \approx 2\xi_u n_q^{\text{eq}}$, etc., with $n_q^{\text{eq}} = 3T^3/\pi^2$, to replace the chemical potentials by particle densities, and we define

$$A(z) = \frac{n_{\tilde{t}}^{\text{eq}}}{n_q^{\text{eq}}} = \frac{z^2}{2}K_2(z). \tag{28}$$

The scatterings induced by $\tilde{u}$ exchange are

$$\begin{aligned} S_{12} - \bar{S}_{12} &= (Y_u + Y_{d_1} + Y_{d_2})\gamma_{S_{12}} + Y_+^{\text{eq}}\gamma_D(B_g\epsilon_{12} + B_{12}\epsilon_g) \\ &\quad -A(z)B_gB_{12}(Y_u + Y_{d_1} + Y_{d_2}) \\ S_{23} - \bar{S}_{23} &= (Y_u + Y_{d_2} + Y_{d_3})\gamma_{S_{23}} + Y_+^{\text{eq}}\gamma_D(B_g\epsilon_{23} + B_{23}\epsilon_g) \\ &\quad -A(z)B_gB_{23}(Y_u + Y_{d_2} + Y_{d_3}) \\ S_{13} - \bar{S}_{13} &= (Y_{d_1} - Y_{d_3})\gamma_{S_{13}} + Y_+^{\text{eq}}\gamma_D(B_{23}\epsilon_{12} - B_{12}\epsilon_{23}) - A(z)B_{12}B_{23}(Y_{d_1} - Y_{d_3}) \\ T_{12} - \bar{T}_{12} &= 2(Y_u + Y_{d_1} + Y_{d_2})\gamma_{T_{12}} \\ T_{23} - \bar{T}_{23} &= 2(Y_u + Y_{d_2} + Y_{d_3})\gamma_{T_{23}} \\ T_{13} - \bar{T}_{13} &= 2(Y_{d_1} - Y_{d_3})\gamma_{T_{13}}. \end{aligned} \tag{29}$$

Here S_{12}, S_{23} and S_{13} correspond to the processes $t\tilde{g} \leftrightarrow \bar{d}_1\bar{d}_2$, $t\tilde{g} \leftrightarrow \bar{d}_2\bar{d}_3$ and $d_1d_2 \leftrightarrow d_2d_3$. Note that for these s-channel processes the resonant parts have been subtracted since they are already included in the Boltzmann equations [13]. T_{12}, T_{23} and T_{13} denote $td_{1,2} \leftrightarrow \tilde{g}\bar{d}_{2,1}$, $td_{2,3} \leftrightarrow \tilde{g}\bar{d}_{3,2}$ and $d_1\bar{d}_{2,3} \leftrightarrow d_{3,2}\bar{d}_2$. Also the exchange of $\tilde{d}_3$ contributes to the above mentioned scatterings

$$\begin{aligned} S'_{12} - \bar{S}'_{12} &= (Y_u + Y_{d_1} + Y_{d_2})\gamma_{S'_{12}} \\ S'_{23} - \bar{S}'_{23} &= (Y_u + Y_{d_2} + Y_{d_3})\gamma_{S'_{23}} \\ T'_{12} - \bar{T}'_{12} &= 2(Y_u + Y_{d_1} + Y_{d_2})\gamma_{T'_{23}} \\ T'_{23} - \bar{T}'_{23} &= 2(Y_u + Y_{d_2} + Y_{d_3})\gamma_{T'_{23}}. \end{aligned} \tag{30}$$

In an approximation we add these contributions incoherently, which is valid when either $\tilde{u}$ or $\tilde{d}_3$ exchange dominates. There are also contributions to the scatterings by exchange of $\tilde{t}$, which we neglect.

For temperatures considerably smaller than the masses of the exchanged squarks we obtain for the scatterings

$$\begin{aligned} \gamma_{S_A} + 2\gamma_{T_A} &= \frac{m_{\tilde{u}}}{H}\frac{352}{3\pi^2}\lambda_{123}^2\alpha_s\frac{1}{z^4}(Y_u + Y_{d_2} + Y_{d_3}) \\ \gamma_{S_R} + 2\gamma_{T_R} &= \frac{m_{\tilde{u}}}{H}\frac{352}{3\pi^2}\lambda_{112}^2\alpha_s\frac{1}{z^4}(Y_u + Y_{d_1} + Y_{d_2}) \\ \gamma_{S_{13}} + 2\gamma_{T_{13}} &= \frac{m_{\tilde{u}}}{H}\frac{352}{16\pi^3}\lambda_{112}^2\lambda_{113}^2\frac{1}{z^4}(Y_{d_1} - Y_{d_3}) \\ \gamma_{S'_A} + 2\gamma_{T'_A} &= \frac{m_{\tilde{u}}}{H}\frac{352}{3\pi^2}\lambda_{123}^2\alpha_s\frac{1}{z^4}\frac{m_{\tilde{u}}^4}{m_{\tilde{d}_3}^4}(Y_u + Y_{d_2} + Y_{d_3}) \\ \gamma_{S'_R} + 2\gamma_{T'_R} &= \frac{m_{\tilde{u}}}{H}\frac{352}{3\pi^2}\theta_{13}^2\lambda_{123}^2\alpha_s\frac{1}{z^4}\frac{m_{\tilde{u}}^4}{m_{\tilde{d}_3}^4}(Y_u + Y_{d_1} + Y_{d_2}). \end{aligned} \tag{31}$$

Note that $\gamma_{S'_A} + 2\gamma_{T'_A}$ is suppressed by a large factor $m_{\tilde{u}}^4/m_{\tilde{d}_3}^4 \sim 1/100$ compared to $\gamma_{S_A} + 2\gamma_{T_A}$ and can be neglected in the following.

Now we can use the relations (25) to eliminate Y_{d_2} and Y_{d_3} in the inverse decay and scattering rates to turn eqs. (20), (27) into a closed set of equations for Y_-, Y_u and Y_{d_1}. The result is given in eq. (44) in the appendix.

4 Analytical and numerical solutions

To understand some features of the numerical solution of the Boltzmann equations (44), let us first discuss some analytical approximations.

The departure from equilibrium in the evolution of Y_+ is of order $1/K$. Since for $m_{\tilde{u}} \lesssim 10^{14}$ GeV we always have $K \gg 1$, we can expand in powers of $1/K$ (see, for instance ref. [18]) To leading order we obtain

$$\begin{aligned} Y_+ - Y_+^{\rm eq} &\approx -\frac{1}{\beta K}\frac{dY_+^{\rm eq}}{dz} \\ &= \frac{1}{Kz}Y_+^{\rm eq} \\ &= \frac{135}{2\pi^4 g_*}\frac{z}{K}K_2(z) \approx \frac{135}{2^{3/2}\pi^{7/2}g_*}\frac{\sqrt{z}}{K}e^{-z}. \end{aligned} \tag{32}$$

Non-equilibrium in the CP-violating densities is governed by some effective value of K, which involves the baryon number violating couplings.

Let us start with the somewhat simpler case $B_{12} = B_{23}$. Then the down quark number densities combine as $Y_{d_1} + 2Y_{d_2} + Y_{d_3}$, since $\gamma_{S_{12},T_{12}} = \gamma_{S_{23},T_{23}}$. (This is not true for the scatterings induced by $\tilde{d}_3$, which however are suppressed by $m_{\tilde{u}}^4/m_{\tilde{d}_3}^4 \sim 1/100$ and can be neglected.) As a result, we can form a closed set of equations for Y_- and Y_B, which is given in the appendix (45). In the limit of large K, the equation for Y_- reduces to

$$Y_- \approx A(z)\left(\frac{5k+2}{8k+2} - \frac{11k+5}{4k+1}B_R\right)Y_B. \tag{33}$$

We also have neglected terms higher order in $A(z)$, which are exponentially small at late times, i.e. for $z \gg 1$. We end up with a single equation for the baryon asymmetry

$$\begin{aligned} \frac{dY_B}{dz} &= \beta K\left\{\frac{\epsilon_g}{2}(Y_+ - Y_+^{\rm eq}) - A(z)B_{12}(1-2B_{12})\frac{11k+5}{4k+1}Y_B\right\} \\ &\quad -\frac{11k+5}{4k+1}Y_B\frac{m_{\tilde{u}}}{H}\frac{352}{3\pi^2}\lambda_{112}^2\alpha_s\frac{1}{z^4}. \end{aligned} \tag{34}$$

We observe that the effective value of K is reduced proportionally to B_{12} and depends somewhat on k.

For large values of $B_{12}K$ we can now solve for Y_B using the method of steepest descent. Assuming that the inverse decays dominate, the freeze-out value of z follows from

$$\frac{e^{z_f}}{z_f^{5/2}} = \frac{\sqrt{\pi}}{2^{3/2}}\frac{11k+5}{4k+1}B_{12}(1-2B_{12})K, \tag{35}$$

which leads to a logarithmic dependence of z_f on K. The final baryon asymmetry is then given by

$$\eta_B = \frac{28}{79}\frac{\epsilon_g}{2}\frac{135}{\pi^{7/2}g_* z_f B_{12}(1-2B_{12})K}\frac{4k+1}{11k+5}\sqrt{\frac{4z_f}{2z_f-5}}, \tag{36}$$

where we added the weak sphaleron factor of 28/79. Using eq. (8) we observe that the leading dependence of η_B on B_{12} cancels. This means that smaller values of the

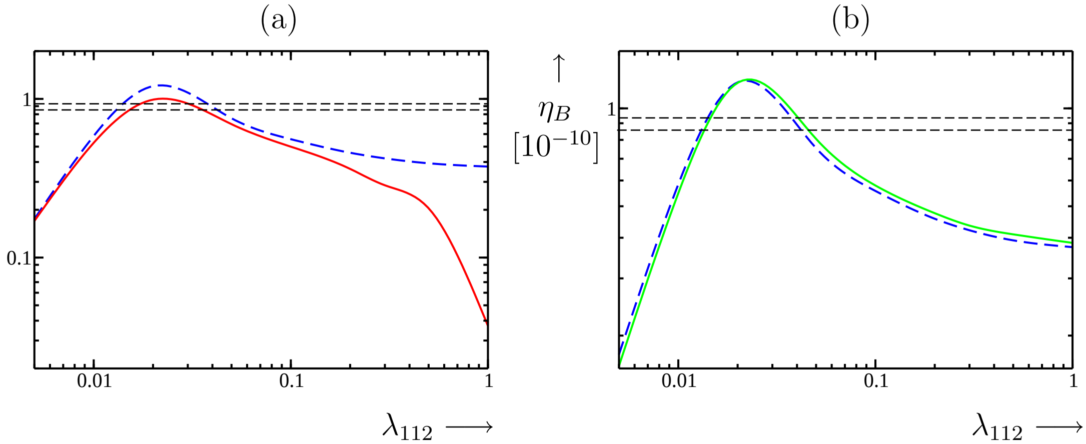

Figure 2: (a) The baryon asymmetry in units of 10^{-10} as a function of λ_{112} for $m_{\tilde{u}} = 5 \times 10^{13}$ GeV, $B_{12} = B_{23}$, $k = 1$, $\theta_{13} = 0.1$ and $m^2_{\tilde{d}_3}/m^2_{\tilde{u}} = 10$. The full (dashed) line includes (neglects) scatterings. The horizontal dashed lines indicate the observed range. (b) The baryon asymmetry with (without) a light $\tilde{q}_1$ in solid (dashed) line. Scatterings are neglected.

baryon number violating couplings reduce by the same amount the CP-asymmetry and the washout by inverse decays. This compensation works until $B_{12}K \sim 1$. Then the washout is no longer important and the baryon asymmetry goes to zero proportional to B_{12}. For very large values of $B_{12}K$ the washout is dominated by scatterings. Then $z_f \propto K^{1/4}$ and the baryon asymmetry is exponentially damped, $\eta_B \propto e^{-4z_f/3}$.

In the general case of $B_{12} \neq B_{23}$ eq. (34) generalizes to

$$\frac{dY_B}{dz} = \beta K \left\{ \frac{\epsilon_{12}}{2}(Y_+ - Y_+^{\rm eq}) - \frac{2}{5}A(z)B_{12}(8 - 7B_{12} - 8B_{23})Y_B \right\}, \qquad (37)$$

where we have neglected scatterings and restricted ourselves to $\epsilon_{23} = 0$ and $k = 1$. Note that the washout term is proportional to B_{12}, i.e. the effective value of K is $K_{\rm eff} \sim B_{12}K$. It can be made small while keeping a large value of B_{23}. This allows us to preserve the d_1 part of the baryon asymmetry in the presence of a large λ_{123}, which can be used to increase the CP asymmetry. A larger value of B_{23} decreases the washout. If ϵ_{23} is non-zero, its contribution to the baryon asymmetry is washed out with $K_{\rm eff} \sim B_{23}K$, i.e. it is not protected by a small value B_{12}.

In the numerical evaluations of eq. (44) we take a maximal CP-phase $\sin\delta = 1$, $\theta_{13} = 0.1$ and $m^2_{\tilde{d}_3}/m^2_{\tilde{u}} = 10$. In fig. 2a we show the baryon asymmetry for $B_{12} = B_{23}$ as a function of λ_{112}. The other parameters are $k = 1$ and $m_{\tilde{u}} = 5 \times 10^{13}$ GeV. The dashed horizontal lines indicate the observed value of the baryon asymmetry, $\eta_B = (0.89 \pm 0.04) \times 10^{-10}$ [19]. The dashed curve takes into account only the inverse decays in the collision terms of eq. (44). In agreement with eq. (36), the baryon asymmetry stays almost constant for $\lambda_{112} \gtrsim 0.02$. For smaller values the

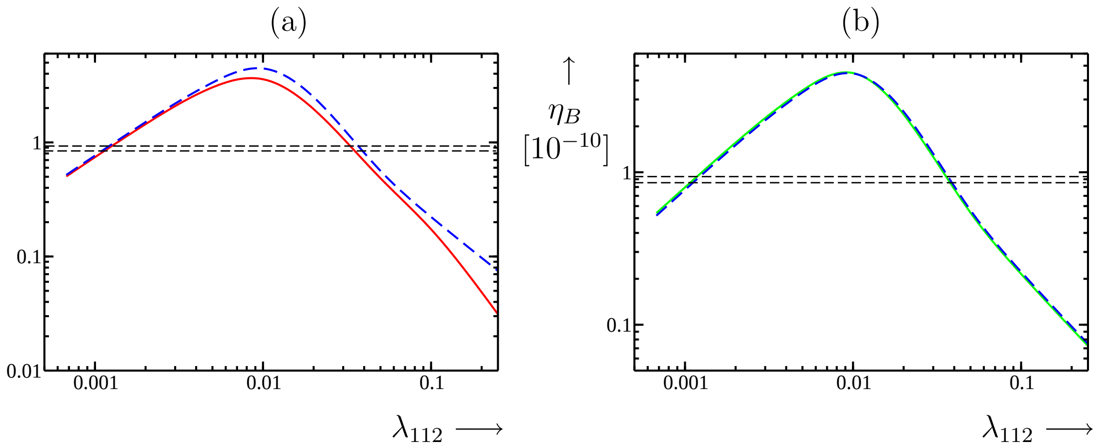

Figure 3: (a) The baryon asymmetry in units of 10^{-10} as function of λ_{112} for $m_{\tilde{u}} = 1 \times 10^{13}$ GeV, $\lambda_{123} = 0.25$, $k_{q_1} = 1$, $\theta_{13} = 0.1$ and $m^2_{\tilde{d}_3}/m^2_{\tilde{u}} = 10$. The full (dashed) line includes (neglects) scatterings. (b) The baryon asymmetry with (without) a light $\tilde{q}_1$ in solid (dashed) line. Scatterings are neglected.

washout becomes negligible and we have $\eta_B \propto B_{12} \propto \lambda^2_{112}$. The solid line takes into account also the scatterings that lead to a strong damping of the baryon asymmetry for $\lambda_{112} \gtrsim 0.2$. In fig. 2b we investigate the influence of a light left-handed squark $\tilde{q}_1$ on the baryon asymmetry, again neglecting scatterings. The difference is quite small since two effects partially cancel: the light $\tilde{q}_1$ ($k = 2$) reduces η_B by adding d.o.f. to the plasma. It also reduces the washout, as we observe from eq. (34). For smaller up squark masses the maximal baryon asymmetry decreases approximately as $\eta_B \propto m_{\tilde{u}}$. So up squark masses above a few times 10^{13} GeV are needed to generate the observed baryon asymmetry in the case $B_{12} = B_{23}$.

A larger baryon asymmetry can be produced if $\lambda_{123} > \lambda_{112}$. In this case ϵ_{12} is enhanced while the washout remains small due to λ_{112}. In fig. 3a we take $\lambda_{123} = 0.25$, $m_{\tilde{u}} = 1 \times 10^{13}$ GeV and $k = 1$. Solid (dashed) lines indicate again that scatterings are included (neglected). Now the baryon asymmetry rises proportionally to $1/\lambda_{112}$. For $\lambda_{112} \lesssim 0.01$ the washout becomes small and $\eta_B \propto \lambda_{112}$. Fig. 3b shows that the inclusion of a light $\tilde{q}_1$ does hardly make any difference. In fig. 4 we compare the baryon asymmetry for $\lambda_{123} = 0.25$ and 0.1. All other parameters are taken as in fig. 3a. We observe that η_B scales approximately as λ_{123}.

The value of λ_{112}, where the baryon asymmetry becomes maximal, i.e. where the washout becomes ineffective, scales as $m_{\tilde{u}}^{-1/2}$. Using this optimal value of λ_{112}, the baryon asymmetry is therefore roughly given by

$$\eta_B \approx 10^{-9} \sin(\delta) \lambda_{123} \frac{\theta_{13}}{0.1} \frac{m^2_{\tilde{u}}/m^2_{\tilde{d}_3}}{0.1} \left(\frac{m_{\tilde{u}}}{10^{13}\text{GeV}} \right)^{1/2} . \tag{38}$$

The observed baryon asymmetry can therefore be generated for $m_{\tilde{u}} \gtrsim 10^{11}$ GeV.[7]

[7]For $m_{\tilde{u}} \lesssim 10^{12}$ GeV the tau and bottom Yukawa interactions as well as the weak sphalerons

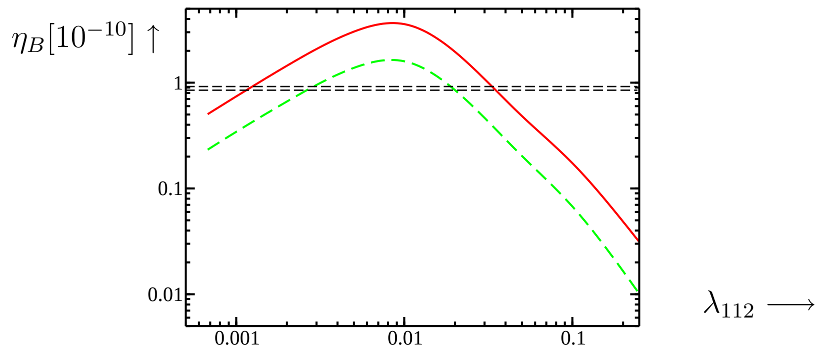

Figure 4: The baryon asymmetry in units of 10^{-10} as function of λ_{112} for $m_{\tilde{u}} = 1 \times 10^{13}$ GeV, $k_{q_1} = 1$, $\theta_{13} = 0.1$ and $m^2_{\tilde{d}_3}/m^2_{\tilde{u}} = 10$. Solid (dashed) lines indicate $\lambda_{123} = 0.25$ (0.1). Scatterings are included.

For $m_{\tilde{u}} = 1 \times 10^{13}$ GeV the coefficient of the off-diagonal gluino vertex can be chosen to be smaller by a factor of 10 than used in the figures.

5 Long-lived neutralinos and gluinos

By exchange of the lightest right-handed squark, $\tilde{u}$, the lightest neutralino predominantly decays as

$$\tilde{\chi}^0_1 \to u d_2 d_3 \tag{39}$$

and the corresponding antiparticles. Again we assumed that $B_{23} \gg B_{12}$, so that the decay into ud_1d_2 is suppressed. The lifetime associated with this process is

$$\tau_{\tilde{\chi}} \approx \frac{1 \times 10^{16}\text{yr}}{y^2 \lambda^2_{123}} \left(\frac{m_{\tilde{\chi}}}{50\text{GeV}}\right)^{-5} \left(\frac{m_{\tilde{u}}}{10^{13}\text{GeV}}\right)^4 . \tag{40}$$

where y denotes the coupling $\chi^0_1 \bar{u}\tilde{u}$. While for a bino-like χ^0_1 one has $y = \frac{2}{3}\sqrt{2}g_1$, this coupling is given by the up Yukawa coupling for a higgsino-like LSP, and even vanishes for a pure wino LSP. Even with $y\lambda_{123} \sim 1$ the LSP lifetime is expected to be much larger than the age of the universe. It can be made even longer lived by suppressing y. If we assume gaugino mass unification, the lightest neutralino is predominantly a bino in the range of small M_2 [4, 20]. Then we estimate $y^2 \sim 1/20$.

Stringent constraints on the lifetime of an instable dark matter particle follow from the production of antiprotons in its decay (39) [21]

$$\tau_{\tilde{\chi}} > 2 \times 10^{19}\text{yr} \left(\frac{m_{\tilde{\chi}}}{50\text{GeV}}\right)^{-1} . \tag{41}$$

can no longer be neglected. We expect the corresponding change in η_B to be at most a factor of order unity.

Similar constraints can also be derived from positron production [22]. This implies that

$$y\lambda_{123} < 0.024 \left(\frac{m_{\tilde{\chi}}}{50\text{GeV}}\right)^{-2} \left(\frac{m_{\tilde{u}}}{10^{13}\text{GeV}}\right)^{2}. \tag{42}$$

It means that for a bino-like LSP we need $m_{\tilde{u}} \gtrsim 10^{13}$ GeV to be in agreement with the anti-proton bound. Light neutralino masses are clearly favored. In split supersymmetry there is the possibility of having a wino- or higgsino as the LSP, with a mass of 2.0-2.5 and 1.0-1.2 TeV, respectively [4, 20]. In this case y can be very much suppressed, but it requires some tuning to keep the bino content at the level of 10^{-3} to compensate for the larger neutralino masses. It seems to be very difficult to use the suppression of y to allow for $m_{\tilde{u}} < 10^{13}$ GeV. Moreover, for very small values of y, loop corrections and the contributions from heavier right-handed squarks become important.

A striking signal of split supersymmetry is the very long-lived gluino. Its lifetime by R-parity conserving decays has been estimated as [23]

$$\tau_{\tilde{g}} = 4\text{sec} \times \left(\frac{m_{\tilde{g}}}{1\text{TeV}}\right)^{-5} \times \left(\frac{\tilde{m}}{10^{9}\text{GeV}}\right)^{4}. \tag{43}$$

Heavy isotope searches induce an upper bound on the sfermion mass scale of $\tilde{m} \lesssim 10^{13}$ GeV [3]. This is consistent with the lower bound we previously derived from the neutralino lifetime. Taking $\tilde{m}(= m_{\tilde{u}}) = 10^{13}$ GeV and $m_{\tilde{g}} = 1$ TeV leads to a gluino lifetime of about 10^{9} yr. If $B_{23} \sim 1$ and $\tilde{u}$ is the lightest squark, the gluino could have a sizable baryon number violating branching ratio.

Constraints from gluino cosmology depend crucially on the gluino annihilation cross section after the QCD phase transition. In ref. [12] it has been argued that the latter is set by the de Broglie wavelength of the gluino rather than the geometric cross section. In this case a stronger bound of $\tilde{m}(= m_{\tilde{u}}) = 10^{12}$ GeV can be derived from diffuse gamma rays. For $m_{\tilde{g}} \gtrsim 500$ GeV an even more restrictive bound is induced by big bang nucleosynthesis. To meet these constraints we would have to assume a mass splitting between the left- and right-handed squarks. The right-handed squarks must be heavier than about 10^{13} GeV to keep the neutralino stable enough. Some left-handed squarks, e.g. $\tilde{q}_1$ have to be lighter than about 10^{12} GeV to speed up gluino decay. In this case the baryon number violating branching ratio of the gluino is highly suppressed by $(\tilde{m}_{q_1}/\tilde{m}_u)^4$. Gauge couplings unification is hardly affected by such a mass splitting. Taking $m_{\tilde{q}_1} = 10^{12}$ GeV, the low energy value of α_s is increased by about 2×10^{-3} with respect to the case where all squarks are degenerate at $\tilde{m} = 10^{13}$ GeV.

6 Conclusions

We have investigated baryogenesis by R-parity violating squark decays in the framework of split supersymmetry. We have restricted ourselves to the baryon number violating couplings λ_{ijk} to avoid rapid neutralino decay. These couplings involve only right-handed squarks. We assume that baryogenesis is dominated by the lightest one, which we take to be the right-handed up squark. A CP-asymmetry in the squark decays arises from the interference of tree-level and 2-loop diagrams involving the CP-phases of λ_{ijk}. The relevant Boltzmann equations include, in addition to the baryon number violating interactions, also the supergauge interactions, the strong sphalerons and the top Yukawa interaction. The generated baryon asymmetry can be enhanced by some hierarchy in the λ_{ijk}, where the CP-asymmetry can be increased while keeping the washout small.

The observed baryon asymmetry can be successfully generated if the up squark mass is larger than 10^{11} GeV. A stronger constraint is induced by neutralino decays. In order to keep the lightest neutralino sufficiently stable to provide the dark matter, the up squark mass has to be at least 10^{13} GeV. Depending on the gluino annihilation after the QCD phase transition, such large squark masses may induce a too large gluino lifetime. However, somewhat lighter left-handed squarks can speed up the gluino decay. At the same time the neutralino remains sufficiently stable and baryogenesis is hardly affected. Favorably, the LSP is bino-like with some higgsino admixture and a mass not far above $M_Z/2$, while the gluino mass is in the few hundred GeV range. Because of the high sfermion mass scale, this scenario predicts sizable corrections to chargino and neutralino Yukawa couplings, which can be probed at a future linear collider [24].

It would be interesting to study also the case of the lepton number violating couplings. These operators could be responsible for the light neutrino masses and one might wonder if this is compatible with successful baryogenesis. However, something would have to be added to the setup to replace the neutralino as the dark matter particle.

Appendix

Using the approximations we discussed in sec. 3, the Boltzmann equations that govern the evolution of Y_-, Y_u and Y_{d_1} take the form

$$\begin{aligned}
\frac{d}{dz}Y_- &= \beta K\left\{-Y_- + A(z)\left[-\left(B_{12}+\frac{8k+4}{5k+2}B_{23}\right)Y_-\right.\right.\\
&\left.\left.+\left(1-2B_{12}-\frac{12k+6}{5k+2}B_{23}\right)Y_u+(B_{23}-B_{12})Y_{d_1}\right]\right\}
\end{aligned}$$

$$\begin{aligned}
\frac{d}{dz}(Y_u-Y_c) &= \frac{2k}{5k+2}\frac{d}{dz}Y_- + \frac{8k+2}{5k+2}\frac{d}{dz}Y_u = \\
&= \beta K\left\{\frac{\epsilon_g}{2}(Y_+-Y_+^{\rm eq}) + B_gY_- - A(z)B_gY_u\right\}\\
&-\frac{m_{\tilde u}}{H}\frac{352}{3\pi^2}\lambda_{123}^2\alpha_s\frac{1}{z^4}\left(1+\frac{m_{\tilde u}^4}{m_{\tilde d_3}^4}\right)\left(\frac{8k+4}{5k+2}Y_- + \frac{12k+6}{5k+2}Y_u - Y_{d_1}\right)\\
&-\frac{m_{\tilde u}}{H}\frac{352}{3\pi^2}\alpha_s\frac{1}{z^4}\left(\lambda_{112}^2+\theta_{13}^2\lambda_{123}^2\frac{m_{\tilde u}^4}{m_{\tilde d_3}^4}\right)(Y_-+2Y_u+Y_{d_1})
\end{aligned}$$

$$\begin{aligned}
\frac{d}{dz}(Y_{d_1}-Y_c) &= \frac{2k}{5k+2}\frac{d}{dz}Y_- + \frac{3k}{5k+2}\frac{d}{dz}Y_u + \frac{d}{dz}Y_{d_1} = \\
&= \beta K\left\{\frac{\epsilon_{12}}{2}(Y_+-Y_+^{\rm eq}) - B_{12}Y_- - A(z)B_{12}(Y_-+Y_u+Y_{d_1})\right\}\\
&-\frac{m_{\tilde u}}{H}\frac{352}{3\pi^2}\alpha_s\frac{1}{z^4}\left(\lambda_{112}^2+\theta_{13}^2\lambda_{123}^2\frac{m_{\tilde u}^4}{m_{\tilde d_3}^4}\right)(Y_-+2Y_u+Y_{d_1})\\
&-\frac{m_{\tilde u}}{H}\frac{352}{16\pi^3}\lambda_{112}^2\lambda_{113}^2\frac{1}{z^4}\left(2Y_{d_1}-\frac{3k+2}{5k+2}Y_- - \frac{2k+2}{5k+2}Y_u\right). \qquad (44)
\end{aligned}$$

The scatterings are given for temperatures considerably smaller than $m_{\tilde u}$, i.e. $z \gg 1$.

In the special case $B_{12} = B_{23}$, the down number densities always appear in the combination $Y_{d_1} + 2Y_{d_2} + Y_{d_3}$, since $\gamma_{S_{12},T_{12}} = \gamma_{S_{23},T_{23}}$. (This is not true for the scatterings induced by $\tilde d_3$, which however are suppressed by $m_{\tilde u}^4/m_{\tilde d_3}^4 \sim 1/100$ and can be neglected.) Thus we can form a closed set of equations for Y_- and Y_B

$$\frac{dY_-}{dz} = \beta K\left\{-Y_- + A(z)\left[\left(-\frac{7k+2}{8k+2}+\frac{7k}{4k+1}B_{12}\right)Y_- + \left(\frac{5k+2}{8k+2}+\frac{11k+5}{4k+1}B_{12}\right)Y_B\right]\right\}$$

$$\begin{aligned}\frac{dY_B}{dz} &= \beta K\left\{\frac{\epsilon_g}{2}(Y_+ - Y_+^{\rm eq}) - 2B_{12}Y_- - A(z)B_{12}\left[\frac{2k}{4k+1}Y_- + \frac{6k+3}{4k+1}Y_B\right]\right\} \\ &\quad - \left(-\frac{5k+2}{4k+1}Y_- + \frac{11k+5}{4k+1}Y_B\right)\frac{m_{\tilde{u}}}{H}\frac{352}{3\pi^2}\lambda_{112}^2\alpha_s\frac{1}{z^4}. \end{aligned} \quad (45)$$

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