Limits on R-parity violation in high-scale supersymmetry

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We revisit the limits on R-parity violation in the minimal supersymmetric standard model. In particular, we focus on the high-scale supersymmetry scenario in which all the sparticles are in excess of the inflationary scale of approximately 10^{13} GeV, and thus no sparticles ever come into thermal equilibrium. In this case, the cosmological limits, stemming from the preservation of the baryon asymmetry that have been previously applied for weak-scale supersymmetry, are now relaxed. We argue that even when sparticles are never in equilibrium, R-parity violation is still constrained via higher dimensional operators by neutrino and nucleon experiments and/or insisting on the preservation of a nonzero $B-L$ asymmetry.

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I. INTRODUCTION

Operators which violate baryon and/or lepton number represent a two-edged sword for beyond the Standard Model (SM) physics. On the one hand, some degree of baryon or lepton number violation is necessary in order to account for the observed baryon asymmetry of the Universe. As is well known, these C- and CP-violating interactions must be out of equilibrium to generate a nonzero asymmetry. An out-of-equilibrium decay, for example, can generate a baryon or lepton asymmetry if the C-, CP-, and B- and/or L-violating decay occurs at a temperature significantly below the mass of the decaying particle [1,2]. A simple rule of thumb condition on the mass, M, of the decaying particle is $M > y^2 M_P$, where y is the coupling leading to the decay and M_P is the (reduced) Planck mass, $M_P^2 = 1/(8\pi G_N)$. In the case that the decay is purely lepton number violating, as in leptogenesis [3], the lepton asymmetry must be converted at least in part to a baryon asymmetry with sphaleron processes [4,5]. However, sphaleron mediated interactions violate $B+L$ while conserving $B-L$ and hence require a nonzero $B-L$ asymmetry to be generated. Most importantly, since the $B-L$ conserving sphaleron processes remain in equilibrium up to the time of the electroweak phase transition, any other process in equilibrium which violates another combination of B and L would lead to the complete washout of any baryon or lepton asymmetry independent of its origin. This allows one to place strong constraints on any possible B- and/or L-violating operators [6–12].

These constraints are particularly important in supersymmetric models with R-parity violation (RPV) [13]. Indeed, R-parity is usually imposed in supersymmetric models to avoid fast baryon and lepton number violating interactions which could lead to rapid proton decay. Limits on proton stability are satisfied if new interactions violated either B *or* L. However, if these interactions remain in equilibrium at the same time as sphalerons are in equilibrium, they would wash away any baryon asymmetry [9,10] despite proton stability. One should bear in mind, that these bounds can be evaded if there is a residual lepton number conservation [14,15], or if the lepton asymmetry is stored in an SM $SU(2)$ singlet such as the right-handed electron [15,16], or other flavor asymmetries [14,17,18]. However, the washout can be affected by slepton mixing angles [19].

Unfortunately, the absence of a supersymmetry signal at the LHC [20] means that the scale of supersymmetry remains unknown. While naturalness can be used to argue for supersymmetry at or near the weak scale [21], the supersymmetry breaking scale may turn out to be much larger. The order parameter for supersymmetry breaking is related to the gravitino mass, $m_{3/2}$, and in minimal anomaly mediated supersymmetry breaking models [22–24], the gravitino mass is typically several hundred TeV to $\mathcal{O}(1)$ PeV. In these models, gaugino masses are loop suppressed with respect to the gravitino mass, and scalar

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masses may be considerably lighter. In models of split supersymmetry [25] and in models of pure gravity mediation [26], the gravitino mass and scalar masses may lie beyond the PeV scale. In models of high-scale supersymmetry [27–29], the gravitino and sparticle masses may be even higher. When the gravitino mass is $\mathcal{O}(1)$ EeV, a new window opens up for gravitino dark matter when all of the sparticle masses (except the gravitino mass) lie above the inflationary scale of approximately 10^{13} GeV [30–32]. However, without some degree of RPV, it is hard to imagine experimental tests to detect EeV gravitino dark matter when all sparticle masses are $> 10^{13}$ GeV.

In high-scale supersymmetry, the limits on RPV are relaxed as the supersymmetric particles were never in the thermal bath and could not participate in interactions that wash out the baryon asymmetry. Therefore, some amount of RPV is acceptable, and if present, RPV operators would render the lightest supersymmetric particle, the gravitino in this case, unstable. If long lived, the decay products may provide a signature for an EeV gravitino [32]. A smoking gun signal could occur from EeV monochromatic neutrinos observable by IceCube and/or ANITA [33,34]. This, in fact is a generic prediction, because given the milder assumptions on RPV couplings, high-scale supersymmetric scenarios are more naturally R-parity violating, allowing more general UV completions where R-parity conservation is not necessary. This compares with weak-scale supersymmetry where RPV couplings need to be significantly suppressed, that either requires an additional suppression mechanism or strongly suggests that R-parity is conserved.

In deriving the limits on RPV parameters, we distinguish between two cases depending on whether or not the gravitino is the dark matter. The limits on RPV parameters are generally stronger when the gravitino lifetime is required to be long enough so as to allow for gravitino dark matter. When these limits are not satisfied, an alternative to gravitino dark matter is required in high-scale supersymmetric models.

In this note, we discuss the cosmological limits on the RPV interactions in high-scale supersymmetry models where all sparticles are assumed to have never been in chemical equilibrium with the SM particles. Thus, we are able to derive new constraints on RPV operators in models of high-scale supersymmetry. The outline of the paper is as follows: in Sec. II, we review previous cosmological constraints on generic higher-dimension operators that arise from the preservation of the baryon asymmetry. In Sec. III, we review and update experimental limits arising from neutrino masses, nucleon decay, and $n-\bar{n}$ oscillations. Specific RPV interactions are discussed in Sec. IV, and we derive new limits for high-scale supersymmetry which are then compared with those from weak-scale supersymmetry. Finally, we also discuss limits that arise from including the gravitino. A summary of our results is given in Sec. V.

II. COSMOLOGICAL LIMITS

Our constraints on RPV interactions are derived from the requirement that $B-L$ violating interactions are not in equilibrium simultaneously with sphaleron processes when they are operative. The sphaleron rate is estimated at next-to-leading order in Refs. [35–37], and sphaleron processes are in thermal equilibrium at temperatures

$$T_c \lesssim T \lesssim T_{\rm sph}, \tag{1}$$

where $T_{\rm sph} \simeq 10^{12}$ GeV and $T_c \simeq 160$ GeV is the critical temperature of the electroweak phase transition [38]. Thus, if there is a $B-L$ violating process in thermal equilibrium at some time when the temperature T satisfies (1), any baryon/lepton number asymmetry is washed out unless it is regenerated after the electroweak phase transition. As noted earlier, there are exceptions to this very general criterion [14–18]. Nevertheless, we are interested in deriving general bounds and to preserve a nonzero asymmetry, we will require all $B-L$ violating processes to be decoupled in the range given by (1). In practice we can distinguish two cases depending on the temperature dependence of the B-/L-violating rate, Γ_{BL}. Assuming that $\Gamma_{BL} \propto T^{2D-7}$ (see below), we require the B-/L-violating rate to be less than the Hubble expansion rate, $\Gamma_{BL} < H$ at $T = T_c$ for $D \leq 4$ (corresponding to relevant or marginal operators) [case (a)], while for $D > 4$ (corresponding to irrelevant operators), it is necessary for $\Gamma_{BL} < H$ at $T = T_{\rm sph}$ [case (b)].

A. Thermal equilibrium constraints

Here, we focus on high-scale supersymmetry scenarios in which all superpartners with the exception of the gravitino have masses greater than the inflationary scale (inflaton mass). As a result, these particles were not produced during reheating and were never in thermal equilibrium. To describe particle interactions in a thermal bath we can use effective operators consisting of only SM particles [9]. These can be written as

$$\mathcal{L}_D = \frac{\mathcal{O}_D}{M_D^{D-4}}, \tag{2}$$

where M_D represents an effective heavy particle mass scale, and $\mathcal{O}_D$ is an operator with mass dimension D. Then, the reaction rate of such an interaction is given by

$$\Gamma_D \sim c_D \left(\frac{T}{M_D}\right)^{2(D-4)} T, \tag{3}$$

with a prefactor c_D which collectively denotes numerical factors arising from the phase space volume and thermal integration for the corresponding reaction. The expressions for c_D are given in Table I. This rate should be compared with the Hubble scale $H = 0.33 g_*^{1/2} T^2/M_P \equiv T^2/\widetilde{M}_P$.

As there are only SM particles in the thermal bath, the number of relativistic degrees of freedom, $g_* = 427/4$, and $\widetilde{M}_P \simeq 7.04 \times 10^{17}$ GeV. Thus, our limits on M_D (which contain all couplings in addition to the heavy mass scale) become

$$M_D > (c_D \widetilde{M}_P T_{\rm c}^{2D-9})^{\frac{1}{2(D-4)}} \quad D < 4, \tag{4}$$

$$M_D > (c_D \widetilde{M}_P T_{\rm sph}^{2D-9})^{\frac{1}{2(D-4)}} \quad D > 4. \tag{5}$$

Next, we consider the reaction rate Γ_D, which can be expressed more accurately as

$$\Gamma_D = \frac{N^2 T}{32\pi^4 n_0} \int_0^\infty ds\, s^{3/2} \sigma(2 \to k) K_1(\sqrt{s}/T), \tag{6}$$

where $\sigma(2 \to k)$ is the $2 \to k$ scattering cross section, K_1 is a modified Bessel function, $n_0 = N\zeta(3)T^3/\pi^2$ is the initial particle number density with $\zeta(3) \simeq 1.202$, and N the number of degrees of freedom for the initial state particle. Note we have neglected the difference between Fermi and Bose statistics, and all of the initial and final state particles are assumed to be massless. The cross sections are proportional to s^{D-5} for the operators with $D \geq 4$, and thus we obtain

$$\int_0^\infty ds\, s^{3/2} s^{D-5} K_1(\sqrt{s}/T) = 2^{2(D-3)} T^{2D-5} \Gamma(D-3)\Gamma(D-2), \tag{7}$$

where Γ is the Gamma function. A concrete form of $\sigma(2 \to k)$ depends on, for example, the spinor and derivative structures of corresponding operators. However, since listing a complete set of $B - L$ breaking operators is not the aim of this paper, we only keep track of a typical phase space volume factor in the following argument. Thus, we may write

$$\sigma(2 \to k) \simeq \frac{1}{2s} |\mathcal{M}(2 \to k)|^2 \Phi_k, \tag{8}$$

where Φ_k is the k-body phase space volume which is approximately given as

$$\Phi_k \simeq \frac{1}{8\pi} \left(\frac{s}{16\pi^2}\right)^{k-2}, \tag{9}$$

with $k \geq 2$.[1]

[1]We implicitly assume the s-channel type decomposition for multiparticle final state diagrams. For more details, see, e.g., Ref. [39].

TABLE I. Cross sections and prefactors in the reaction rates are shown for each operator with effective mass scale M_D. Although we do not consider dimension six operators when applying cosmological limits, as they do not violate $B - L$, we show them in the table for completeness. For $D = 3$, 4, we denote λ as a generic coupling constant and M as the mass of the decaying particle. We have neglected factors of N related to the number of degrees of freedom of incoming states and used Maxwell-Boltzmann statistics in the thermal average.

D	Operator	$\Gamma(1 \to 2)$		c_D
3	ϕ^3	$\frac{\lambda^2}{8\pi} M$		$\frac{\lambda^2}{16\zeta(3)\pi}$
4	$\phi\psi^2$	$\frac{\lambda^2}{8\pi} M$		$\frac{\lambda^2}{16\zeta(3)\pi}$
D	Operator	$\sigma(2 \to k)$		c_D
4	ϕ^4	$\frac{1}{16\pi}\frac{\lambda^2}{s}$	$(k = 2)$	$\frac{\lambda^2}{128\zeta(3)\pi^3}$
5	$\psi^2\phi^2$	$\frac{1}{16\pi}\frac{1}{M_5^2}$	$(k = 2)$	$\frac{1}{16\zeta(3)\pi^3}$
6	ψ^4	$\frac{1}{16\pi}\frac{s}{M_6^4}$	$(k = 2)$	$\frac{3}{2\zeta(3)\pi^3}$
7	$\psi^4\phi$	$\frac{1}{16\pi(4\pi)^2}\frac{s^2}{M_7^6}$	$(k = 3)$	$\frac{9}{2\zeta(3)\pi^5}$
	$\psi^4\mathcal{D}$	$\frac{1}{16\pi}\frac{s^2}{M_7^6}$	$(k = 2)$	$\frac{72}{\zeta(3)\pi^3}$
9	ψ^6	$\frac{1}{16\pi(4\pi)^4}\frac{s^4}{M_9^{10}}$	$(k = 4)$	$\frac{2700}{\zeta(3)\pi^7}$
10	$\psi^6\phi$	$\frac{1}{16\pi(4\pi)^6}\frac{s^5}{M_{10}^{12}}$	$(k = 5)$	$\frac{28350}{\zeta(3)\pi^9}$

For the $D = 3, 4$ operators, one-to-two processes are also possible. When we parametrize the decay width for a particle with mass M to two Dirac fermions by

$$\Gamma(1 \to 2) \equiv \frac{\lambda^2}{8\pi} M, \tag{10}$$

with a generic coupling λ, the reaction rate for $D = 3$, 4 becomes

$$\Gamma_D \simeq \frac{N}{2\pi^2 n_0} M T^2 \Gamma(1 \to 2) = \frac{\lambda^2}{16\zeta(3)\pi} \frac{M^2}{T}, \tag{11}$$

and thus we have c_3, $c_4(\text{decay}) = \lambda^2/(16\zeta(3)\pi)$. For a two-to-two process with quartic coupling λ, we find

$$\Gamma_D \simeq \frac{\lambda^2 N}{128\zeta(3)\pi^3} T, \tag{12}$$

where $c_4(\text{scattering}) = \lambda^2 N/(128\zeta(3)\pi^3)$ for two-to-two processes.[2]

Table I summarizes the decay width and cross sections along with the coefficients c_D which are used to limit M_D for each operator, and where $\psi, \phi, \mathcal{D}$ symbolically denote a fermion, scalar, and covariant derivative, respectively.

[2]This result differs from that in [10] by a factor of $6/\pi^2(2)$ for the decay (two-to-two) processes that result from our approximation of using Maxwell-Boltzmann statistics in the thermal average.

TABLE II. Cosmological and laboratory limits on baryon and/or lepton number violating operators. The cosmological limits assume the SM value of $g_* = 427/4$ as appropriate for high-scale supersymmetry.

				Cosmological limits		
D	Operator	ΔB	ΔL	case (a)	case (b)	Laboratory limits
5	$llhh$	0	2	$M_5 > 10^{8.6}$ GeV	$M_5 > 10^{13.5}$ GeV	$M_5 > 10^{14.7}$ GeV
6	$qqql$	1	1	none	none	$M_6 > 10^{15.8}$ GeV
7	$\psi^4 h(qqql^c h, \text{etc.})$	$\not{B}-\not{L}$		$M_7 > 10^{4.5}$ GeV	$M_7 > 10^{12.6}$ GeV	$M_7 > 10^{11.1}$ GeV
	$\psi^4 \mathcal{D}(qqql^c \mathcal{D}, \text{etc.})$	$\not{B}-\not{L}$		$M_7 > 10^{4.8}$ GeV	$M_7 > 10^{13.0}$ GeV	$M_7 > 10^{10.1}$ GeV
9	$qqqqqq$	2	0	$M_9 > 10^{3.8}$ GeV	$M_9 > 10^{12.6}$ GeV	$M_9 > 10^{5.9}$ GeV
	$qqqlll$	1	3	$M_9 > 10^{3.8}$ GeV	$M_9 > 10^{12.6}$ GeV	$M_9 > 10^{5.6}$ GeV
10	$qqql^c l^c l^c h$	1	-3	$M_{10} > 10^{3.5}$ GeV	$M_{10} > 10^{12.5}$ GeV	$M_{10} > 10^{5.1}$ GeV

Again, in the table we omit prefactors in the cross sections coming from the kinematics and coupling structure. For dimension seven operators involving a covariant derivative, we consider only derivative terms, since they are dominant compared to the terms with a gauge boson. The resultant limits on the operators are summarized in Table II where $q, l(l^c), h$ symbolically represent quark, lepton (charge conjugate of a lepton field), and the SM Higgs fields, respectively.

III. LABORATORY LIMITS

In addition to the cosmological limits discussed above, there are of course a variety of laboratory limits on baryon and lepton number violating operators which we summarize in this section. We will concentrate on limits from neutrino masses (on lepton number violating operators) and nucleon decay limits.

A. Neutrino mass constraints

We can derive a lower bound on the mass scale M_5 used in the dimension five (Weinberg) operator, $(L^T C \cdot H)\times(L \cdot H)/M_5$, in addition to the cosmological bound in Eq. (5), from the upper bound on the sum over the neutrino masses, $\sum_i m_{\nu_i} < 0.15$ eV (95% CL) [40]. In addition, we can derive an upper bound on M_5 from neutrino oscillation data. To be more concrete, in the following discussion we assume that the heaviest neutrino mass is given by $m_\nu = v^2/M_5$ where $v = 1/\sqrt{2\sqrt{2}G_F} \simeq 174$ GeV is the vacuum expectation value of the SM Higgs field.[3]

Since the sign of the squared neutrino mass difference cannot be determined from atmospheric neutrino observations, the neutrino mass ordering can either be the normal hierarchy (NH) or the inverted hierarchy (IH). The NH spectrum is defined by the neutrino mass ordering given by $m_{\nu_3} > m_{\nu_2} > m_{\nu_1}$, whereas the IH spectrum is given by $m_{\nu_2} > m_{\nu_1} > m_{\nu_3}$. The central values for the experimentally determined mixing angles and squared mass differences are [41][4]

$$\begin{aligned}
&\sin^2\theta_{12} = 0.307, \quad \sin^2\theta_{23} = 0.542(\text{NH}),\\
&\qquad 0.536(\text{IH}), \quad \sin^2\theta_{13} = 2.18\times 10^{-2},\\
&m_S^2 \equiv m_{\nu_2}^2 - m_{\nu_1}^2 = 7.53\times 10^{-5}\ \text{eV}^2,\\
&m_A^2 \equiv |m_{\nu_3}^2 - m_{\nu_2}^2| = 2.53\times 10^{-3}\ \text{eV}^2(\text{NH}),\\
&\qquad 2.44\times 10^{-3}\ \text{eV}^2(\text{IH}).
\end{aligned} \tag{13}$$

For the neutrino masses, there is only one free parameter in each case, namely, the lightest neutrino mass which we denote $m_{\nu 0}$. Then, each mass spectrum is given by

$$\begin{aligned}
\text{NH:}\ & m_{\nu_1} = m_{\nu 0}, \quad m_{\nu_2} = \sqrt{m_{\nu 0}^2 + m_S^2},\\
& m_{\nu_3} = \sqrt{m_A^2 + m_S^2 + m_{\nu 0}^2},
\end{aligned} \tag{14}$$

$$\begin{aligned}
\text{IH:}\ & m_{\nu_1} = \sqrt{m_A^2 - m_S^2 + m_{\nu 0}^2},\\
& m_{\nu_2} = \sqrt{m_A^2 + m_{\nu 0}^2}, \quad m_{\nu_3} = m_{\nu 0}.
\end{aligned} \tag{15}$$

The limit $\sum_i m_{\nu_i} < 0.15$ eV sets an upper bound, $m_{\nu 0} \lesssim 4.16\times 10^{-2}$ eV (3.24×10^{-2} eV) corresponding to the heaviest neutrino mass, $m_{\nu_3(\nu_2)} \simeq 6.59\times 10^{-2}$ eV (5.91×10^{-2} eV) for the NH (IH) case. In both cases, we obtain $M_5 \gtrsim 10^{14.7}$ GeV. This constraint is listed in the first line of Table II.

On the other hand, neutrino oscillations imply nonzero neutrino masses which gives an upper bound to M_5, assuming that the only source for generating neutrino masses is the dimension five operator. To obtain a conservative limit, we take $m_{\nu 0} = 0$ in both cases, and then the heaviest mass becomes $m_{\nu_3(\nu_2)} \simeq 5.10\times 10^{-2}$ eV (4.94×10^{-2} eV) for the

[3]Note that the Yukawa coupling, y, has been absorbed into the definition of M_5, so that if the Weinberg operator arises from a right-handed neutrino mass and the see-saw mechanism then $M_5 = M_R/y^2$.

[4]For $\sin^2\theta_{23}$, we are using fits in the second octant.

NH (IH) case. It turns out that in both cases we need $M_5 \lesssim 10^{14.8}$ GeV to explain the neutrino oscillation data, provided the dimension five operator is the dominant contribution to the neutrino masses.[5] The combined upper and lower limits point to a unique value in the range of $10^{14.7}$ GeV $\lesssim M_5 \lesssim 10^{14.8}$ GeV, though one needs to bear in mind that M_5 is an effective mass parameter for the heaviest neutrino that includes all relevant couplings.

B. Nucleon decays

Clearly, most baryon number violating operators can be constrained by proton or neutron decay. Since the dimension six operator $qqql$ does not violate $B-L$, there is no cosmological limit from the preservation of the baryon asymmetry. On the other hand, as baryon number is violated, the nucleon lifetime can be used to derive a lower bound on M_6.

In general, there are four types of such operators [42–44],[6]

$$\mathcal{O}^{(1)}_{ijkl} = G^{(1)}_{ijkl}(\bar{d}^c_{i\alpha}P_R C\bar{u}^{cT}_{j\beta})(Q^T_{k\gamma}C\cdot L_l)\epsilon^{\alpha\beta\gamma}, \quad (16)$$

$$\mathcal{O}^{(2)}_{ijkl} = G^{(2)}_{ijkl}(Q^T_{i\alpha}C\cdot Q_{j\beta})(\bar{u}^c_{k\gamma}P_R C\bar{e}^{cT}_{l})\epsilon^{\alpha\beta\gamma}, \quad (17)$$

$$\mathcal{O}^{(3)}_{ijkl} = G^{(3)}_{ijkl}(Q^T_{i\alpha}C\cdot Q_{j\beta})(Q^T_{k\gamma}C\cdot L_l)\epsilon^{\alpha\beta\gamma}, \quad (18)$$

$$\mathcal{O}^{(4)}_{ijkl} = G^{(4)}_{ijkl}(\bar{d}^c_{i\alpha}P_R C\bar{u}^{cT}_{j\beta})(\bar{u}^c_{k\gamma}P_R C\bar{e}^{cT}_{l})\epsilon^{\alpha\beta\gamma}, \quad (19)$$

where i, j, k, l and α, β, γ are flavor and color indices, respectively, and the $SU(2)$ product is denoted as $A\cdot B = \epsilon^{ab}A_{aL}B_{bL}$ $(\epsilon^{12}=\epsilon^{123}=1)$, i.e., $Q^T C\cdot L = u^T C P_L e - d^T C P_L \nu$. All fermions are defined as four component spinors, and the $SU(2)$ doublets include the chiral projection P_L as appropriate. The flavor dependent Wilson coefficients are represented by $G^{(n)}_{ijkl}$ $(n=1,2,3,4)$, whose flavor structure depends on the underlying theory. Instead of specifying a concrete flavor structure, we assume that there are neither any degenerate parameters (i.e., no accidental cancellations) nor large hierarchies among the different flavor entries.

For proton decay involving a charged antilepton in the final state, the strongest limit arises from the decay mode $p\to e^+\pi^0$, which is given by $\tau_{p\to e\pi} > 1.6\times 10^{34}$ years [46]. The relevant operators for this decay channel are

$$\begin{aligned} &G^{(1)}_{1111}(\bar{d}^c P_R C\bar{u}^{cT})(u^T C P_L e),\\ &-2G^{(2)}_{1111}(d^T C P_L u)(\bar{u}^c P_R C\bar{e}^{cT}),\\ &-2G^{(3)}_{1111}(d^T C P_L u)(u^T C P_L e),\\ &G^{(4)}_{1111}(\bar{d}^c P_R C\bar{u}^{cT})(\bar{u}^c P_R C\bar{e}^{cT}), \end{aligned} \quad (20)$$

where the flavor mixing from the CKM matrix is neglected. The decay width of this process is given by

$$\Gamma(p\to e^+\pi^0) = \frac{m_p}{32\pi}\left(1-\frac{m_\pi^2}{m_p^2}\right)^2(|\mathcal{A}_L|^2+|\mathcal{A}_R|^2), \quad (21)$$

where $m_p(m_\pi)$ is the proton (pion) mass. The amplitudes are defined as[7]

$$\mathcal{A}_L = G^{(1)}_{1111}\langle\pi^0|(ud)_R u_L|p\rangle - 2G^{(3)}_{1111}\langle\pi^0|(ud)_L u_L|p\rangle, \quad (22)$$

$$\mathcal{A}_R = -2G^{(2)}_{1111}\langle\pi^0|(ud)_L u_R|p\rangle + G^{(4)}_{1111}\langle\pi^0|(ud)_R u_R|p\rangle, \quad (23)$$

where we have used the notation

$$(ud)_\Gamma u_{\Gamma'} = \epsilon^{\alpha\beta\gamma}(u^T_\alpha C P_\Gamma d_\beta)P_{\Gamma'}u_\gamma, \quad (24)$$

with $\Gamma,\Gamma' = L$, R, and α, β, γ are $SU(3)$ color indices [48]. Then, by taking $G^{(1,2,3,4)}_{1111}\sim 1/M_6^2$ and the hadron matrix elements $\langle\pi^0|(ud)_{L,R}u_{L,R}|p\rangle \sim -\langle\pi^0|(ud)_{L,R}u_{R,L}|p\rangle \sim 0.1$ GeV^2 [48], we obtain $M_6 > 10^{15.8}$ GeV. This constraint is listed in the second line of Table II. If some of the dimension six operators are absent, this constraint is somewhat relaxed. For instance, when $\mathcal{O}^{(1)}$ is the sole operator mediating proton decay, the limit reduces to $M_6 > 10^{15.5}$ GeV.

A similar procedure can be applied for the operators whose mass dimension is greater than six.

(1) Dimension seven operators: The $qqql^c h$ type operators may be regarded as a dimension six operator where the Wilson coefficient, estimated as $\sim v/M_7^3$, leads to the limit $M_7 \gtrsim 10^{11.1}$ GeV. In the same way, the $qqql^c\mathcal{D}$ type operators may be regarded as dimension six operators with the Wilson coefficient $\sim\Lambda_{\rm QCD}/M_7^3$ with $\Lambda_{\rm QCD}\simeq 300$ MeV [41], and thus we obtain $M_7 \gtrsim 10^{10.1}$ GeV. It should be noted that for dimension seven operators, there are additional operators that violate $B-L$, such as $(e^{cT}CL\cdot H)\times(L^T C\cdot L)$.[8] It is also known that all dimension seven

[5]If there is another source for the neutrino masses and the dimension five operator is not dominant, then this upper limit does not apply, and M_5 can be as large as possible, e.g., M_P.

[6]Our definition of these operators is identical to that of Ref. [45] after arranging $SU(2)$ and spin indices appropriately.

[7]Note that in the specific case of SU(5), proton decay mediated by X and Y gauge bosons only involves $G^{(1)}$ and $G^{(2)}$ [47].

[8]There also exist operators involving more than two Higgs doublets or covariant derivatives (equivalently, gauge field strength), which we do not discuss here.

operators that violate either baryon or lepton number, violate $B-L$ as well. For more on all such operators, see Refs. [49–51].

(2) Dimension nine operators: The $qqqlll$ operator may also induce nucleon decays. For instance, the operator $(L^T C\cdot L)(e^{cT}CP_L d^c)(\bar{Q}\cdot C\bar{Q}^T)$ causes $n\to e^+e^-\nu$ which has the lifetime constraint $\tau_{n\to e^+e^-\nu} > 2.8\times 10^{32}$ years [52]. Then, we obtain $M_9 > 10^{5.6}$ GeV where we have estimated the decay width as $\Gamma(n\to e^+e^-\nu)\simeq \alpha_h^2 m_n^5/(256\pi^3 M_9^{10})$ (by only taking into account the phase space volume) with the hadron matrix element $\alpha_h \simeq -0.0144\ \mathrm{GeV}^3$ [53].[9] The $qqqqqq$ type of operators, especially $uddudd$, are constrained by $n-\bar{n}$ oscillation. Following Refs. [57,58], the $n-\bar{n}$ mixing time can be written as $\tau_{n-\bar{n}}\simeq 1/\delta m$ with

$$\delta m \sim \frac{1}{M_9^5}|\psi(0)|^2, \tag{25}$$

where $\psi(0)$ denotes the neutron wave function at the origin, which is typically $\psi(0)\sim \Lambda_{\mathrm{QCD}}^3$. The current constraint $\tau_{n-\bar{n}} > 2.7\times 10^8$ s [59] sets the limit $M_9 > 10^{5.9}$ GeV.

(3) Dimension ten operators: For instance, $(L^T C\cdot L)(e^{cT}CP_L d^c)(\bar{Q}\cdot Hu^c)$ induces the nucleon decay $n\to e^+e^-\nu$ whose decay width may be evaluated as $\Gamma_{n\to e^+e^-\nu}\simeq \alpha_h^2 m_n^5 v^2/(256\pi^3 M_{10}^{12})$, and thus we obtain $M_{10} > 10^{5.1}$ GeV.

Finally let us reiterate that for operators of mass dimension higher than seven there exist many baryon and/or lepton number violating operators which are not listed in Table II. Though the $D=7$, 9, 10 operators in Table II are just examples, it is sufficient for our purpose since such higher dimensional operators would usually involve more undetermined parameters compared to lower dimensional ones, and thus the detailed constraints are strongly model dependent. For instance, some operators that violate lepton number, but conserve baryon number, can be constrained by neutrinoless double beta decay [60], while the limit strongly depends on the form of the operators. Nevertheless, when we assume that baryon and lepton numbers are violated at the same scale, and that there is no large hierarchy between the mass scales of baryon and lepton number violation, the constraints on M_D from nucleon decays are in most cases stronger than those from neutrinoless double beta decay (see, e.g., Ref. [61].)

[9]This constraint was missing in Ref. [54], and an accurate estimate has recently been given in Ref. [55] which also argues that $n\to K^0 l^+ l^-\nu$, induced by, for example, $(L^T C\cdot L)\times(e^{cT}CP_L d^c)(u^{cT}CP_L d^c)$, is larger than $n\to l^+l^-\nu$ in RPV models, although this particular channel is not yet constrained by experiments (see also Ref. [56]).

IV. *R*-PARITY VIOLATING INTERACTIONS

We now discuss the limits on RPV interactions using the results obtained in the previous sections. The RPV superpotential is given by

$$W_{\mathrm{RPV}} = W_{\mathrm{RPV}}^{(2)} + W_{\mathrm{RPV}}^{(3)}, \tag{26}$$

$$W_{\mathrm{RPV}}^{(2)} = \mu_i' H_u\cdot L_i, \tag{27}$$

$$W_{\mathrm{RPV}}^{(3)} = \frac{1}{2}\lambda_{ijk}L_i\cdot L_j E_k^c + \lambda_{ijk}' L_i\cdot Q_i D_k^c + \frac{1}{2}\lambda_{ijk}'' U_i^c D_j^c D_k^c. \tag{28}$$

The explicit Lagrangian including soft supersymmetry breaking terms is shown in the Appendix. We will first review the bounds derived in the case of weak-scale supersymmetry [10] and contrast them with bounds obtained in high-scale supersymmetry. These bounds are derived from both the cosmological preservation of the baryon asymmetry and the experimental limits on baryon and/or lepton number violating processes including proton decay. We will also comment on the limits on the RPV parameters when we require a sufficiently long-lived gravitino as the dark matter.

In general, the RPV mass parameter μ_i' depends on lepton flavor, but here we omit the flavor dependence for simplicity and take $\mu_i'\equiv\mu'$. (For a more detailed discussion, see, e.g., [62,63].) Since lepton number is not conserved, L and H_d cannot be distinguished, and thus there is a field basis dependence in defining L and H_d fields. For instance, if $L\to(1-\epsilon^2)^{1/2}L+\epsilon H_d$ and $H_d\to(1-\epsilon^2)^{1/2}H_d-\epsilon L$ with $\epsilon=\mu'/\sqrt{\mu^2+\mu'^2}$ and μ is the μ-parameter in the minimal supersymmetric SM (MSSM) superpotential, we can eliminate the bilinear RPV term at the expense of generating trilinear RPV terms, such as $y_u\epsilon LLE^c$ and $y_d\epsilon QLD^c$. For simplicity and since observables do not depend on the choice of basis, we will work in the basis that explicitly keeps the bilinear term (27) given in W_{RPV}.

A. Limits on μ'

1. Weak-scale supersymmetry

As discussed above, there are strong constraints on baryon and lepton number violating operators whose induced interactions are simultaneously in equilibrium with the sphaleron interactions. In the case of an R-parity violating bilinear LH_u term, one-to-two processes involving a Higgsino, lepton, and a gauge boson will be induced. From Eq. (11), the thermally averaged rate at a temperature, T for these lepton number violating interactions is given by [10,32]

$$\Gamma_{1\to 2} = \frac{g^2\theta^2 T}{16\zeta(3)\pi} \simeq 0.016 g^2 \frac{\mu'^2}{m_f^2} T, \tag{29}$$

where g is a gauge coupling, and $\theta \simeq \mu'/m_f$ is the mixing angle induced by μ' for a fermion with mass m_f. We require that this lepton number violating interaction is out of equilibrium. As such, we require the interaction rate (29) is less than the Hubble rate, $H \simeq \sqrt{\pi^2 g_*/90} T^2/M_P$. This implies that

$$\mu'^2 < 20\sqrt{g_*}\frac{T^3}{M_P}, \tag{30}$$

where the fermions have a thermal mass, $m_f \sim gT$. We further insist that any lepton number violating rate involving μ' remains out of equilibrium while sphaleron interactions are in equilibrium, i.e., between the weak scale T_c and $T_{\rm sph}$. As one can see, the limit (30) is strongest for T of order the weak scale [case (a) corresponding to $D = 3$]. For weak-scale supersymmetry, the fermion can be either a lepton or Higgsino, $g_* = 915/4$ and at T_c one obtains the limit [10]

$$\mu' < 2.3 \times 10^{-5}\ \text{GeV}. \tag{31}$$

For weak-scale supersymmetry, this limit translates to $\epsilon \lesssim 2.3 \times 10^{-7}$.

In general, the RPV bilinear term induces a nonzero neutrino mass via a dimension five operator. The mixing angle between neutrinos and the Higgsino is given by μ'/μ, and through the Higgsino-Higgs-gaugino (wino or bino) coupling, we obtain a dimension five operator of the form

$$\mathcal{L}_5 \simeq \frac{1}{M_5}\nu_L\nu_L hh, \qquad \frac{1}{M_5} \simeq \epsilon^2 \frac{g_2^2 M_1 + g_1^2 M_2}{M_1 M_2(1+\tan^2\beta)}, \tag{32}$$

where $M_1(M_2)$ are the bino (wino) masses and $g_2(g_1)$ is the $SU(2)_L(U(1)_Y)$ gauge coupling.

In weak-scale supersymmetry models, the limit (31) is stronger than the limit from neutrino masses [62,64] which comes from the dimension five operator with the constraint given in Table II. As one can see from Table II, the strongest limit from a dimension five two-to-two process is obtained by requiring the out-of-equilibrium condition to hold at the highest possible scale, which in this case is $T_{\rm sph}$ [case (b)]. For weak-scale supersymmetry, the limit on M_5 becomes

$$M_5 > \frac{(c_5 T M_P)^{1/2}}{\sqrt{0.33 g_*^{1/2}}} \approx 2.8 \times 10^{13}\ \text{GeV}, \tag{33}$$

for $T = T_{\rm sph}$ and $g_* = 915/4$ (the change in g_* accounts for the slight difference with respect to the limit in Table II). This translates to the limit

$$\mu' < 1.9 \times 10^{-7}\ \text{GeV}^{-1/2}\tilde{m}^{1/2}\mu(1+\tan^2\beta)^{1/2}/g \approx 4.4 \times 10^{-4}\ \text{GeV}, \tag{34}$$

for $\mu \sim M_1 \sim M_2 \sim \tilde{m} \sim 100$ GeV, and $\tan\beta \approx 1$. We assume a generic gauge coupling $g \sim 0.6$ throughout. In this case, $\epsilon \lesssim 4.4 \times 10^{-6}$.

2. High-scale supersymmetry

In the case of high-scale supersymmetry, we assume that all sparticles are heavier than the inflationary mass scale $m_I \sim 3 \times 10^{13}$ GeV, and we denote the typical sparticle mass scale as $\tilde{m} > m_I$. As all sparticle masses are greater than $T_{\rm sph}$, there are no sparticles in the thermal bath when sphalerons are in equilibrium and the limit from one-to-two processes is not applicable. Nevertheless, the limit from the effective dimension five operator is valid when the heavy sparticles are integrated out. Since only Standard Model particles are in the thermal bath, $g_* = 427/4$ and we can use the limit on M_5 from Table II [case (b)]. The limit on μ' becomes

$$\mu' < 1.7 \times 10^{-7}\ \text{GeV}^{-1/2}\tilde{m}^{1/2}\mu(1+\tan^2\beta)^{1/2}/g \approx 6.6 \times 10^{13}\ \text{GeV}, \tag{35}$$

for $\mu \sim M_1 \sim M_2 \sim \tilde{m} \sim 3 \times 10^{13}$ GeV, and $\tan\beta \approx 1$. In this case, $\epsilon \lesssim 2.2$.

As one can also see from Table II, the laboratory limit in this case is in fact the strongest limit on μ'. Using $M_5 > 5 \times 10^{14}$ GeV, we obtain

$$\mu' < 4.5 \times 10^{-8}\ \text{GeV}^{-1/2}\tilde{m}^{1/2}\mu(1+\tan^2\beta)^{1/2}/g \approx 1.7 \times 10^{13}\ \text{GeV}, \tag{36}$$

or $\epsilon \lesssim 0.57$.

Note that if $W^{(2)}_{\rm RPV}$ is the only source of neutrino mass, our previous limit on $M_5 < 10^{14.8}$ GeV translates into a lower bound on μ',

$$\mu' > 4 \times 10^{-8}\ \text{GeV}^{1/2}\tilde{m}^{1/2}\mu(1+\tan^2\beta)^{1/2}/g \approx 1.5 \times 10^{13}\ \text{GeV}. \tag{37}$$

As discussed in Sec. III, the lower limit can be removed if there is another source for generating neutrino masses that can explain the neutrino oscillation data.

B. Limits on $\lambda,\lambda',\lambda''$

1. $D = 4, 5$

The quartic couplings in Eq. (28) can lead to either one-to-two processes (involving a scalar and two fermions) or two-to-two processes (involving four scalars) which violate baryon and/or lepton number. The rates for these processes taken from Table I can be written as [10]

$$\Gamma_{2\to 2} = \frac{\lambda^2 y^2 T}{128\zeta(3)\pi^3} \simeq 2\times 10^{-4}\lambda^2 y^2 T, \tag{38}$$

$$\Gamma_{1\to 2} = \frac{\lambda^2 m_0^2}{16\zeta(3)\pi T} \simeq 0.016\lambda^2 \frac{m_0^2}{T}, \tag{39}$$

where λ is a generic RPV quartic coupling in (28) and $m_0 < T$ is the scalar mass. The rate (38) depends on the Standard Model Yukawa coupling y, because the baryon/lepton number violating processes actually arise from a cross term in the F-term in the scalar potential.

In weak-scale supersymmetry, these processes will be in equilibrium unless λ is quite small, and the limit on λ is derived by comparing these rates with the Hubble rate. This yields the limits

$$\lambda < 1.2\times 10^{-6} y^{-1} \quad 2\leftrightarrow 2, \tag{40}$$

$$\lambda < 1.4\times 10^{-7} \quad 1\leftrightarrow 2, \tag{41}$$

where we have evaluated the limit at $T\sim m_0 \sim T_c$ in the one-to-two rate.

Once again, in the case of high-scale supersymmetry, when all sparticle masses are greater than the inflationary scale, the above limits are no longer applicable as there are no sparticles in the thermal bath at the time when sphaleron interactions are in equilibrium. For the RPV bilinear term, we were able to derive a limit on μ' by integrating out the heavy sparticles and setting a limit on the resulting dimension five operator. One might think that one can do the same for the quartic coupling, and form a dimension six (four-fermion) operator and still set (weaker) limits on the RPV quartic couplings. However, as shown in [42], there are no $B-L$ violating dimension six operators involving only Standard Model fields.

There are, however, numerous laboratory and astrophysical constraints on the RPV quartic couplings which are independent of the sphaleron processes [65–71]. For example, some of the quartic couplings will contribute radiatively to neutrino masses and neutrinoless double beta decay [65], where these limits scale as $\lambda < \mathcal{O}(10^{-3})(\tilde{m}/100\ \mathrm{GeV})^p$ for $p = 1/2, 5/2$, respectively. As one can see, in the high-scale supersymmetric limit, the bounds on these couplings also disappear. The same is true of collider limits [68] and cosmological and astrophysical limits from the decay of the lightest supersymmetric particle [69].

Furthermore, let us comment on the issue of radiatively induced neutrino masses [72]. Possible radiative corrections through the RPV couplings are summarized in Ref. [73], where the relevant contribution in our case is the self-energy diagrams (diagram 19 in that paper) involving μ' and $B'(\equiv B_i)$. The correction to the neutrino mass, δm_ν, is proportional to $\mu' B'/\tilde{m}^2$. However, once all the Higgs boson contributions are incorporated, one finds that δm_ν is suppressed by $(v/\tilde{m})^2$, and thus δm_ν may be written in terms of the dimension five operator with a loop factor, i.e., $(16\pi^2)^{-1}(L^T C\cdot H)(L\cdot H)/M_5$ with $1/M_5 \sim \mu' B'/\tilde{m}^4 \sim \mu'^2/\tilde{m}^3$. Therefore, the constraint from δm_ν is weaker than that coming from the tree level (dimension five) operator.

2. D = 6

Despite the weakening of most bounds on the RPV quartic couplings, there remain limits on dimension six operators which induce proton decay.[10] Once again, since these operators conserve $B-L$, there are no limits from the sphaleron washout of the baryon asymmetry. Nevertheless, proton decay is induced by $\tilde{d}^c$ exchange diagrams in the RPV case [42], and only $\mathcal{O}^{(1)}$ type of operators can appear. The corresponding Wilson coefficients are

$$G^{(1)}_{ijkl} \simeq \sum_{m,n=1}^{3} \lambda''^*_{jim}\lambda'_{lkn}(m^{-2}_{\tilde{d}^c})_{mn}, \tag{42}$$

where the relation $\lambda''_{ijk} = -\lambda''_{ikj}$ is imposed by gauge symmetry. By ignoring flavor mixing in the down-type squark sector, we obtain

$$\mathcal{O}^{(1)}_{1111} = \frac{1}{M_6^2}(\bar{d}^c_1 P_R C \bar{u}^{cT}_1)(Q_1^T C\cdot L_1),$$
$$\frac{1}{M_6^2} \sim \frac{1}{\tilde{m}^2}\sum_{m=1}^{3}\lambda''^*_{11m}\lambda'_{11m}, \tag{43}$$

and thus the limit on M_6 from Table II can be expressed as the following limit on the quartic coupling:

$$\left|\sum_{m=1}^{3}\lambda''^*_{11m}\lambda'_{11m}\right| < 2.3\times 10^{-5}\left(\frac{\tilde{m}}{3\times 10^{13}\ \mathrm{GeV}}\right)^2, \tag{44}$$

which updates the results given in Ref. [62].

3. D = 7

Dimension seven operators of the type, $qqql^c h$, are induced by involving trilinear couplings, i.e., $A^d_{ij} h_d\cdot \tilde{Q}_i \tilde{d}^c_j$ and $\mu y^{d*}_{ij} h_u^\dagger \cdot \tilde{Q}_i \tilde{d}^c_j$. Then, we have

$$\mathcal{L}^{(1)}_7 = G^{(1)}_{7,ijkl}(\bar{L}_i\cdot h_d d_j)(Q_l^T C\cdot Q_k), \tag{45}$$

$$\mathcal{L}^{(2)}_7 = G^{(2)}_{7,ijkl}(\bar{L}_i\cdot h_u^\dagger d_j)(Q_l^T C\cdot Q_k), \tag{46}$$

with coefficients

$$G^{(1)}_{7,ijkl} \simeq \lambda'^*_{imj}(m^{-2}_{\tilde{Q}})_{mm'}A^d_{m'n'}(m^{-2}_{\tilde{d}^c})_{n'n}\lambda''^*_{kln}, \tag{47}$$

[10]See [74] for related discussion on proton decay constraints.

$$G^{(2)}_{7,ijkl} \simeq \lambda'^*_{imj}(m^{-2}_{\tilde{Q}})_{mm'}\mu^* y^d_{m'n'}(m^{-2}_{\tilde{d}^c})_{n'n}\lambda''^*_{kln}, \quad (48)$$

respectively, where μ is assumed to be complex. These operators give rise to interaction rates which scale as

$$\Gamma_7 \simeq c_7(\lambda'\lambda'')^2\left(\frac{A^2}{\tilde{m}^8}\right)T^7, \quad (49)$$

with $c_7 = 9/2\zeta(3)\pi^5$ given in Table I, and A denotes an A term. When compared to the Hubble rate, one sees that the appropriate limit, evaluated at $T = T_{\rm sph}$ [case (b)], gives

$$\left|\sum_{m=1}^{3}\lambda'^*_{imj}\lambda''^*_{klm}\right| < 290\left(\frac{\tilde{m}}{3\times 10^{13}\ \text{GeV}}\right)^4\left(\frac{3\times 10^{13}\ \text{GeV}}{|A^d_0|}\right), \quad (50)$$

$$\left|\sum_{m,n=1}^{3}\lambda'^*_{imj}y^d_{mn}\lambda'^*_{klm}\right| < 290\left(\frac{\tilde{m}}{3\times 10^{13}\ \text{GeV}}\right)^4\left(\frac{3\times 10^{13}\ \text{GeV}}{|\mu|}\right), \quad (51)$$

where $(m^2_{\tilde{Q}})_{ij} \sim (m^2_{\tilde{d}^c})_{ij} \sim \tilde{m}^2\delta_{ij}$ and $A^d_{ij} = A^d_0\delta_{ij}$, and we have assumed that there is no flavor mixing in the soft supersymmetry breaking terms. These cosmological limits are stronger than the nucleon decay limits which are

$$\left|\sum_{m=1}^{3}\lambda'^*_{1m1}\lambda''^*_{11m}\right| < 10^7\left(\frac{\tilde{m}}{3\times 10^{13}\ \text{GeV}}\right)^4\left(\frac{3\times 10^{13}\ \text{GeV}}{|A^d_0|}\right), \quad (52)$$

$$\left|\sum_{m,n=1}^{3}\lambda'^*_{1m1}y^d_{mn}\lambda'^*_{11m}\right| < 10^7\left(\frac{\tilde{m}}{3\times 10^{13}\ \text{GeV}}\right)^4\left(\frac{3\times 10^{13}\ \text{GeV}}{|\mu|}\right), \quad (53)$$

and become very weak in the high-scale supersymmetric limit.

Through the trilinear couplings $A^u_{ij}h_u\cdot\tilde{Q}_i\tilde{u}^c_j$ and $\mu y^{u*}_{ij}h^\dagger_d\cdot\tilde{Q}_i\tilde{u}^c_j$, we also have

$$\mathcal{L}^{(3)}_7 = G^{(3)}_{7,ijkl}(\bar{L}_i\cdot h_u C\bar{d}^{cT}_j)(\bar{d}^c_l C\bar{d}^{cT}_k), \quad (54)$$

$$\mathcal{L}^{(4)}_7 = G^{(4)}_{7,ijkl}(\bar{L}_i\cdot h^\dagger_d C\bar{d}^{cT}_j)(\bar{d}^c_l C\bar{d}^{cT}_k), \quad (55)$$

with coefficients

$$G^{(3)}_{7,ijkl} \simeq \lambda'^*_{imj}(m^{-2}_{\tilde{Q}})_{mm'}A^u_{m'n'}(m^{-2}_{\tilde{u}^c})_{n'n}\lambda''^*_{kln}, \quad (56)$$

$$G^{(4)}_{7,ijkl} \simeq \lambda'^*_{imj}(m^{-2}_{\tilde{Q}})_{mm'}\mu^* y^u_{m'n'}(m^{-2}_{\tilde{u}^c})_{n'n}\lambda''^*_{kln}, \quad (57)$$

and the constraints on $G^{(3)}_{7,ijkl}$ and $G^{(4)}_{7,ijkl}$ can be obtained in the same way.

4. D = 9, 10

RPV interactions also induce $n-\bar{n}$ oscillations via dimension nine operators, which can be written in the following form [71]:

$$\mathcal{L}_9 \supset G_{9,ijklmn}(\bar{d}^c_i P_R C\bar{u}^{cT}_j)(\bar{d}^c_k P_R C\bar{u}^{cT}_l)(\bar{d}^c_m P_R C\bar{d}^{cT}_n). \quad (58)$$

There are two possible diagrams that produce this operator, namely, via the A-term or gluino exchange. In each case, we obtain

$$G_{9,ijklmn}(A-\text{term}) \simeq \sum_{ss'tt'uu'}\frac{\lambda''^*_{jis}\lambda''^*_{lkt}\lambda''^*_{unm}A''_{s't'u'}}{(m^2_{\tilde{u}^c})_{ss'}(m^2_{\tilde{d}^c})_{tt'}(m^2_{\tilde{d}^c})_{uu'}}, \quad (59)$$

$$G_{9,ijklmn}(\text{gluino}) \simeq \sum_{m'n'}\frac{g_s}{M_3}\frac{\lambda''^*_{ijm'}\lambda''^*_{lkn'}}{(m^2_{\tilde{d}^c})_{m'm}(m^2_{\tilde{d}^c})_{n'n}}, \quad (60)$$

where M_3 is the gluino mass, g_s is the QCD coupling, and A'' is a soft mass parameter (see Appendix). The rate for these processes can be approximated as

$$\Gamma_9 = c_9\lambda''^6\left(\frac{A''^2}{\tilde{m}^{12}}\right)T^{11}, \quad A-\text{term}, \quad (61)$$

$$\Gamma_9 = c_9\lambda''^4 g_s^2\left(\frac{1}{\tilde{m}^{10}}\right)T^{11}, \quad \text{gluino}, \quad (62)$$

with $c_9 = 2700/\zeta(3)\pi^7$ given in Table I. Then, once again comparing to the Hubble rate at $T = T_{\rm sph}$, we obtain the constraints as follows:

$$\left|\sum_{stu}\lambda''^*_{jis}\lambda''^*_{lkt}\lambda''^*_{unm}\left(\frac{A''_{stu}}{3\times 10^{13}\ \text{GeV}}\right)\right| < 3.4\times 10^4\left(\frac{\tilde{m}}{3\times 10^{13}\ \text{GeV}}\right)^6, \quad (63)$$

$$|\lambda''^*_{ijm}\lambda''^*_{lkn}| < 5.6\times 10^4\left(\frac{\tilde{m}}{3\times 10^{13}\ \text{GeV}}\right)^5, \quad (64)$$

where $M_3 \sim \tilde{m}$ and $(m^2_{\tilde{u}^c})_{ij} \simeq (m^2_{\tilde{d}^c})_{ij} \sim \tilde{m}^2\delta_{ij}$. It is also true that the above case (b) limit leads to a stronger bound than the $n-\bar{n}$ oscillation limits which are effectively absent,

$$\left|\sum_{stu}\lambda''^*_{11s}\lambda''^*_{11t}\lambda''^*_{u11}\left(\frac{A''_{stu}}{3\times 10^{13}\ \text{GeV}}\right)\right| < 7.7\times 10^{37}\left(\frac{\tilde{m}}{3\times 10^{13}\ \text{GeV}}\right)^6, \quad (65)$$

$$|\lambda''^*_{111}\lambda''^*_{111}| < 1.3\times 10^{38}\left(\frac{\tilde{m}}{3\times 10^{13}\ \text{GeV}}\right)^5. \quad (66)$$

We also have dimension nine operators of the type $qqqlll$, which can be expressed as

$$\mathcal{L}_9 \supset G_{9,ijklmn}(L_i^T C \cdot L_j)(e_k^{cT} C P_L d_l^c)(u_m^{c\,T} C P_L d_n^c), \quad (67)$$

$$G_{9,ijklmn} \simeq \sum_{ss'tt'uu'} \lambda_{sjk}\lambda'_{itl}\lambda''_{mnu}(m_{\tilde{L}}^{-2})_{ss'}(m_{\tilde{Q}}^{-2})_{tt'}(m_{\tilde{d}^c}^{-2})_{uu'}A'^{*}_{s't'u'}. \quad (68)$$

The decay $n \to Kl\bar{l}\nu$ occurs through this operator, although this particular channel has not been experimentally constrained. The decay $n \to l\bar{l}\nu$ happens through $(L^T C_i \cdot L_j)(e_k^{cT} C P_L d_l^c)(\bar{Q}_m \cdot C\bar{Q}_n^T)$ in dimension nine. However, as discussed in Ref. [55], such operators should be loop suppressed. This may be understood by comparing with the operator $(L_i^T C \cdot L_j)(e_k^{cT} C P_L d_l^c)(u_m^{c\,T} C P_L d_n^c)$ in which one of the d^c quarks should be an s quark because of $SU(3)$ color symmetry. A chiral flip is needed to avoid the appearance of s quark. For instance, once the diagram involves a Higgs boson loop, $(u_m^{c\,T} C P_L d_n^c)$ may be replaced by $y^u y^d/16\pi^2 \times (\bar{Q}_m \cdot C\bar{Q}_n^T)$ which is less restrictive for the RPV couplings.

We may also construct a dimension ten operator of the type $qqql^c l^c l^c h$ by looking at dimension nine operators. For example, d_n^c in $(L_i^T C \cdot L_j)(e_k^{cT} C P_L d_l^c)(u_m^{c\,T} C P_L d_n^c)$ may be replaced by $m_d/v \times (C\bar{Q}_n^T \cdot H)$ to yield $(L^T C_i \cdot L_j)\times(e_k^{cT} C P_L d_l^c)(\bar{Q}_n \cdot H u_m^c)$. However, for these higher dimensional operators, the cosmological limit is much stronger than the laboratory limits as seen in the dimension seven operator case.

C. Limits from the gravitino

The limits on μ' in high-scale supersymmetry from Sec. IV A 2 are based solely on the preservation of the baryon asymmetry and experimental limits on the neutrino mass. However, if we insist that the lightest supersymmetric particle, the gravitino, is relatively stable so that it can play the role of dark matter, we can derive a significantly stronger limit on μ' [32]. The presence of the RPV parameter μ' opens up channels for gravitino decay into neutrinos plus gauge/Higgs bosons. The total decay rate is [32]

$$\Gamma_{3/2} \simeq \frac{\epsilon^2 \cos^2\beta m_{3/2}^3}{16\pi M_P^2}, \quad (69)$$

where $\epsilon \approx \mu'/\mu \approx \mu'/m_I$ and $\cos\beta \approx 1/\sqrt{2}$. Demanding that the gravitino lifetime exceeds the current age of the Universe ($\tau_{3/2} > 4.3\times10^{17}$ s) corresponds to a limit on $\mu' < 0.03$ GeV, for $m_I = 3\times10^{13}$ GeV and $m_{3/2} = 1$ EeV. Note that the upper limit on μ' scales as $m_I/m_{3/2}^{3/2}$. An even more restrictive limit on μ' is possible from the IceCube constraints on the neutrino flux produced by the gravitino decay [33]. In this case, we require a lifetime $\tau_{3/2} > 10^{28}$ s which corresponds to a limit, $\mu' < 2\times10^{-7}$ GeV. As one can see, these limits are far more restrictive than those from baryon/lepton number conservation and neutrino masses.

While the quartic RPV couplings can also induce gravitino decay, they do so only at the one-loop level. As a consequence, the limit on a generic quartic coupling is very weak. For example, writing a generic dimension six operator as $(fff\psi_\mu)/M_6^2$ with f and ψ_μ being SM fermions and the gravitino, respectively, the corresponding cosmological limit on M_6 from the preservation of $B-L$ asymmetry is

$$M_6 > (c_6 \tilde{M}_P T_{\rm sph}^3)^{1/4} \simeq 1.3\times10^{13}\ \text{GeV}, \quad (70)$$

with $c_6 = 3/2\zeta(3)\pi^3$ given in Table I. The same dimension six operator makes the gravitino unstable, and the gravitino decay width into three SM fermions is estimated as

$$\Gamma_{3/2\to3f} \simeq \frac{m_{3/2}^5}{256\pi^3 M_6^4}. \quad (71)$$

Assuming the gravitino lifetime is longer than the age of the Universe gives rise to the lower bound

$$M_6 > 5.3\times10^{20}\ \text{GeV}\left(\frac{m_{3/2}}{\text{EeV}}\right)^{5/4} \quad (\tau_{3/2} > 4.3\times10^{17}\ \text{s}). \quad (72)$$

Furthermore, the operator $(\bar{d}^c P_R C \bar{u}^{cT})(\bar{d}^c \gamma^\mu \psi_\mu)$ can be induced at the one-loop level, with a suppression scale $M_6^{-2} \sim |\lambda''|^2 A''/(16\pi^2\tilde{m}^2 M_P) \sim |\lambda''|^2/(16\pi^2\tilde{m}M_P)$. From the cosmological limit (70), this gives

$$|\lambda''| \lesssim 8.2\times10^3, \quad (73)$$

when $\tilde{m} \sim m_I = 3\times10^{13}$ GeV. Again imposing the gravitino lifetime limit (72), we obtain

$$|\lambda''| \lesssim 2\times10^{-4}\left(\frac{\text{EeV}}{m_{3/2}}\right)^{5/4}. \quad (74)$$

In this case, the limit on the quartic coupling is only competitive with the limit in (40) when the gravitino mass begins to approach the inflationary scale, or when $m_{3/2} \gtrsim 0.002\times m_I$. If instead we impose the IceCube gravitino lifetime limit, $\tau_{3/2} > 10^{28}$ s [33], the generic constraint becomes

$$M_6 > 2.1\times10^{23}\ \text{GeV}\left(\frac{m_{3/2}}{\text{EeV}}\right)^{5/4} \quad (\tau_{3/2} > 10^{28}\ \text{s}), \quad (75)$$

and thus we obtain

$$|\lambda''| \lesssim 5 \times 10^{-7} \left(\frac{\mathrm{EeV}}{m_{3/2}}\right)^{5/4}. \tag{76}$$

This limit is more stringent than (40) when $m_{3/2} \gtrsim 1.7 \times 10^{-5} m_I$.

Finally note that there may be additional RPV operators which are nonrenormalizable corrections to the superpotential and the Kähler potential. These operators can also contribute to the baryon number violating interactions. For example, consider the operator in Eq. (2.72) of [62] with coupling κ_7. By introducing a supersymmetry breaking spurion superfield X (containing the Goldstino ψ), we may write the corresponding term as

$$K \supset \frac{\kappa_7}{M_P^2} Q \cdot Q d^{c\dagger} X^\dagger, \tag{77}$$

which gives $(\kappa_7/M_P^2)(\bar{d}^c Q) \cdot (\bar{\psi} Q)$. Then by identifying $M_6 = M_P/\sqrt{\kappa_7}$, the cosmological limit becomes $\kappa_7 \lesssim 3.4 \times 10^{10}$. This will lead to a dimension nine operator induced by gravitino exchange, which remains below the Hubble rate so long as $\kappa_7 < 2.1 \times 10^{11}$.[11] However, this operator could also lead to gravitino decay and using Eq. (71), we can derive a much stronger limit, $\kappa_7 < 2 \times 10^{-5} (\mathrm{EeV}/m_{3/2})^{5/2}$ for $\tau_{3/2} > 4.3 \times 10^{17}$ s, and $\kappa_7 < 1.4 \times 10^{-10} (\mathrm{EeV}/m_{3/2})^{5/2}$ for $\tau_{3/2} > 10^{28}$ s. We do not consider these nonrenormalizable operators any further.

V. SUMMARY

We have revisited the limits on RPV in the minimal supersymmetric standard model, where the constraints from laboratory experiments and the preservation of the $B - L$ asymmetry are discussed. In particular, we have focused on the high-scale supersymmetry scenario in which all the sparticles are in excess of the inflationary scale of approximately 10^{13} GeV, and thus they were never in equilibrium. Since the previously argued cosmological limits in weak-scale supersymmetry assume that the sparticles involved in the RPV interactions remain in equilibrium at T_c, those limits cannot be applied to the high-scale supersymmetry case, and thus, the cosmological limits from the preservation of a nonzero $B - L$ asymmetry are relaxed. Nevertheless, even when sparticles are never in equilibrium, RPV couplings are still constrained through higher dimensional operators.

Based on effective operators, we have reviewed and updated the experimental and cosmological limits on B- and/or L-violating processes, and then applied them for RPV in the high-scale supersymmetry scenario. For dimension five operators, we have shown that the neutrino mass constraints are stronger than the cosmological limit, while for operators of mass dimension higher than seven, the cosmological limit is stronger than the experimental limits. Dimension six operators are only constrained by nucleon decay experiments since there are no $B - L$ violating operators of dimension six. We have contrasted the limits on RPV in high-scale supersymmetry with those in weak-scale supersymmetry up to dimension ten operators and shown that indeed a wider range of RPV couplings is acceptable. This implies that unlike weak-scale supersymmetry, high-scale supersymmetry can generically have RPV with mild constraints on the couplings and imposing an R-parity is not necessary. This then leads to the generic prediction of an EeV gravitino decay.

We have also distinguished limits based on the assumption of gravitino dark matter. In this case, the RPV interactions lead to the possibility of gravitino decay. If long lived, the gravitino may still provide a sufficient mass density to make up the dark matter. Indeed, if very long lived, present day decays may yet provide a signature. For example, the RPV bilinear proportional to μ' induces a decay to neutrinos which could be seen in high energy neutrino detectors [32]. If we require the presence of dark matter today, or sufficiently long lived so as not to surpass the existing experimental constraints, we obtain limits on μ' which are far stronger than those from baryon/lepton number violation. The RPV quartic couplings on the other hand are better constrained by baryon/lepton number violating rates. We also noted that our limits can be applied to nonrenormalizable corrections in supergravity models, with the most stringent limits arising from gravitino decay.

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APPENDIX: NOTATIONS AND LAGRANGIAN

We summarize the relevant Lagrangian we have used in our discussion for the sake of completeness. Our notations and conventions basically follow Appendixes A and B of Ref. [62].[12] In this appendix, we recall some useful

[11] Here we assume $m_{3/2} < T_{\mathrm{sph}}$ and estimate the reaction rate as $\kappa_7^4 T^9/M_P^8$ which should be compared with H at $T = T_{\mathrm{sph}}$.

[12] We use different labeling for some fields, but there is an obvious correspondence with Ref. [62]. For clarity, we also denote the $SU(2)$ products with a dot.

relations when writing the operators in terms of $SU(2)_L$ doublet fields, and then the relevant parts in the Lagrangian are presented.

Four-component Dirac spinors for leptons are constructed as

$$e = \begin{pmatrix} \psi_e \\ \bar{\psi}_{e^c} \end{pmatrix}, \qquad \nu = \begin{pmatrix} \psi_\nu \\ \bar{\psi}_\nu \end{pmatrix}, \tag{A1}$$

where the two-component Weyl spinor ψ_i denotes the corresponding fermionic part in superfields. Then, we may write the $SU(2)_L$ product for the doublet L as

$$L_i^T C \cdot L_j = \nu_i^T C P_L e_j - e_i^T C P_L \nu_j, \tag{A2}$$

with i and j being the flavor indices. We also note that in this notation $P_L e^c$ refers to the corresponding right-handed field, i.e., ψ_{e^c}. We may define four component spinors for quarks in the same manner.

The MSSM superpotential is given by

$$W_{\rm MSSM} = \mu H_u \cdot H_d + y_{ij}^e H_d \cdot L_i E_j^c + y_{ij}^d H_d \cdot Q_i D_j^c + y_{ij}^u H_u \cdot Q_i U_j^c, \tag{A3}$$

where $y^{e,u,d}$ are the Yukawa coupling matrices. The RPV trilinear couplings are given by

$$W_{\rm RPV}^{(3)} = \frac{1}{2}\lambda_{ijk} L_i \cdot L_j E_k^c + \lambda'_{ijk} L_i \cdot Q_j D_k^c + \frac{1}{2}\lambda''_{ijk} U_i^c D_j^c D_k^c, \tag{A4}$$

and the corresponding Lagrangian in terms of four-component fermions becomes

$$\mathcal{L}_{L_i L_j E_k^c} = -\frac{1}{2}\lambda_{ijk}[L_i^T C \cdot L_j \tilde{e}_k^c + e_k^{cT} C \tilde{L}_i \cdot L_j] + \text{H.c.}, \tag{A5}$$

$$\mathcal{L}_{L_i Q_j D_k^c} = -\lambda'_{ijk}[\tilde{d}_k^c Q_j^T C \cdot L_j + d_k^{cT} C \tilde{L}_i \cdot Q_j + d_k^{cT} C \tilde{Q}_j \cdot L_i] + \text{H.c.}, \tag{A6}$$

$$\mathcal{L}_{U_i^c D_j^c D_k^c} = -\frac{1}{2}\lambda''_{ijk}[\tilde{u}_i^c (d_j^{cT} C P_L d_k^c) + \tilde{d}_j^c (u_i^{cT} C P_L d_k^c) + \tilde{d}_k^c (u_i^{cT} C P_L d_j^c)] + \text{H.c.}, \tag{A7}$$

where Q, L are quark, lepton doublets satisfying $Q^T C \cdot L = u^T C P_L e - d^T C P_L \nu$, and $\lambda_{ijk} = -\lambda_{jik}$ and $\lambda''_{ijk} = -\lambda''_{jik}$ due to the gauge symmetry.

The soft supersymmetry breaking terms consist of the RPV part,

$$-\mathcal{L}_{\rm RPV}^{\rm soft} = \frac{1}{2}A_{ijk}\tilde{L}_i \cdot \tilde{L}_j \tilde{e}_k^c + A'_{ijk}\tilde{L}_i \cdot \tilde{Q}_j \tilde{d}_k^c + \frac{1}{2}A''_{ijk}\tilde{u}_i^c \tilde{d}_j^c \tilde{d}_k^c + B_i h_u \cdot \tilde{L}_i + \tilde{m}_{di}^2 h_d^\dagger \tilde{L}_i + \text{H.c.}, \tag{A8}$$

while the R-parity conserved part is given by

$$\begin{aligned} -\mathcal{L}^{\rm soft} = {} & (m_{\tilde{Q}}^2)_{ij}\tilde{Q}_i^\dagger \tilde{Q}_j + (m_{\tilde{u}^c}^2)_{ij}\tilde{u}_i^{c\dagger}\tilde{u}_j^c + (m_{\tilde{d}^c}^2)_{ij}\tilde{d}_i^{c\dagger}\tilde{d}_j^c + (m_{\tilde{L}}^2)_{ij}\tilde{L}_i^\dagger \tilde{L}_j + (m_{\tilde{e}^c}^2)_{ij}\tilde{e}_i^{c\dagger}\tilde{e}_j^c + (A_{ij}^e h_d \cdot \tilde{L}_i \tilde{e}_j^c + A_{ij}^d h_d \cdot \tilde{Q}_i \tilde{d}_j^c \\ & - A_{ij}^u h_u \cdot \tilde{Q}_i \tilde{u}_j^c + \text{H.c.}) + \tilde{m}_d^2 h_d^\dagger h_d + \tilde{m}_u^2 h_u^\dagger h_u + (B h_u \cdot h_d + \text{H.c.}) + \frac{1}{2}M_1 \bar{\tilde{B}}\tilde{B} + \frac{1}{2}M_2 \bar{\tilde{W}}^3 \tilde{W}^3 + M_2 \bar{\tilde{W}}^+ \tilde{W}^+ \\ & + \frac{1}{2}M_3 \bar{\tilde{g}}^a \tilde{g}^a, \end{aligned} \tag{A9}$$

where $\tilde{Q}$ and $\tilde{L}$ are squark and slepton doublets, respectively.

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