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COSMOLOGICAL CONSTRAINTS ON B AND L VIOLATION

BY

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The undersigned certify that they have read, and recommend to the Faculty of Graduate Studies and Research for acceptance, a thesis entitled "Cosmological Constraints on B and L Violation" submitted by Sacha Davidson in partial fulfillment of the requirements for the degree of Doctor of Philosophy in Theoretical Physics.

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ABSTRACT

The existence of the baryon asymmetry of the Universe puts strong constraints on extensions of the Standard Model which violate baryon and/or lepton number. Interactions violating baryon number (B) but conserving lepton number (L) in the early Universe could wash away any previously established baryon asymmetry. Interactions which violate lepton number separately, with or without associated violation of baryon number, could combine with non-perturbative electroweak effects to eradicate the cosmological baryon asymmetry. We derive constraints on such interactions arising from the persistence of the cosmological baryon asymmetry. After implementing astrophysical constraints, we discuss the prospects for observing B and/or L violation in laboratory experiments. Modulo loopholes that we mention, we find that even if R-parity is violated, the lifetime of the lightest supersymmetric particle must be so long that its decays could not be observed in accelerator experiments, and that L-violating Z decays would have unobservably small branching ratios. The only novel signature for accelerator experiments that survives our analysis is a small window for the lightest supersymmetric particle to be strongly-interacting or charged, with a lifetime that is short on a cosmological time-scale but long enough to appear stable in accelerator experiments. We also find that if $\Delta B = 2$ interactions exist, the rates they yield for $n - \bar{n}$ oscillations and $N - N$ annihilations in nuclei are probably below the present observational threshold.

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To all the inhabitants, past and present, of Silmarillion Toggenburgs:
in particular,

to Feanar, for keeping me sane,

to Nicklebloom, because she could have done better physics than I,

to all the "Top Ten" does, especially Taurwen; if I had worked as hard as you, this thesis would be twice as long,

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INTRODUCTION

1.1 Introduction

It is a misfortune of present day particle physics that everything we can calculate within the Standard Model seems to agree with experiment [1]. This is wonderful for the model, but boring for physicists, and not altogether in line with our prejudices about what a theory should look like. We therefore assume that there is physics beyond the Standard Model (SM), and compute constraints on new theories by requiring that they agree with observation.

It appears that non-perturbative electroweak processes that violate $B + L$ (which I will frequently call "sphalerons", although this is not accurate), but conserve $B - L$, become strong enough at high temperatures to be in thermal equilibrium in the early Universe above $T_c \gtrsim 100$ GeV [2, 3]. Any previously produced $B + L$ asymmetry would therefore be taken to its equilibrium value. Unless the Baryon Asymmetry of the Universe (BAU) is produced after the "sphalerons" drop out of equilibrium ($T \equiv T_c$), there must be an initial asymmetry in $B - L$ to preserve $B \neq 0$. Furthermore, for the asymmetry to survive at $T > T_c$, no other interaction violating B or L may simultaneously be in thermal equilibrium. This requirement puts strong upper bounds on the effective coupling constants of interactions violating any combination of B and L other than that taken to zero by the "sphalerons" [4, 5, 6, 7, 8, 9] (modulo extra symmetries and mass effects ...see chapter 4).

This introduction consists of a short review of the Standard Model [10], supersymmetry [11], and the connection between instantons, anomalies and $B + L$ violation in the early Universe. There is no review of cosmology (see [12]). The second chapter is a quick overview of the argument constraining B, L and $B-L$ violating operators in which estimates of the upper bounds on the interaction rates are calculated. This material was originally presented in reference [7]: "Cosmological baryon asymmetry constraints on extensions of the Standard Model". In the third chapter we calculate the constraints on B or L violating operators more carefully, and look at the implications of those constraints for laboratory experiments (This comes from reference [8]). The fourth chapter is a more detailed look at what the equilibrium asymmetries between particle and antiparticle densities are, and at the implications of extra global symmetries for our constraints. This was motivated by a recent preprint of Ibanez and Quevedo [13] pointing out that our bounds on higher dimensional supersymmetric operators are too optimistic if there are extra global symmetries.

1.2 The Standard Model

a) Gauge theories

One of the most familiar global internal symmetries is that of phase invariance; one may chose the overall phase of a complex field at will. However, it seems intuitively peculiar that one must make the same transformation at all points in space-time—in a fuzzy way one might expect to be able to chose different origins for the coordinates of an internal symmetry at different points (particularily if these are separated by a space-like distance).

If a Lagrangian $\mathcal{L}(\phi, \partial_\mu \phi)$ is invariant under the global transformation

$$\phi \longrightarrow e^{i\theta}\phi \tag{1.1}$$

then one can make the transformation local by making θ space-time dependent. This, however, is no longer a symmetry of $\mathcal{L}$ because the derivative of ϕ picks up an extra term:

$$\partial_\mu \phi \longrightarrow \partial_\mu \phi + i\phi\partial_\mu\theta. \tag{1.2}$$

So one introduces a new field with an unphysical longitudinal degree of freedom that can be used to remove the unwanted terms in the kinetic energy. (In a discretized model of space-time, A_μ is a phase degree of freedom living on the links between lattice sites [14].) The new ("covariant") derivative is

$$D_\mu \equiv \partial_\mu - iA_\mu \tag{1.3}$$

and the transformation under phase rotations of the vector field is

$$A_\mu \longrightarrow A_\mu - i\partial_\mu\theta \tag{1.4}$$

The kinetic term for the new field is

$$K.T. = -\frac{1}{4g^2}F_{\mu\nu}F^{\mu\nu} \tag{1.5}$$

where g is the gauge coupling and

$$F_{\mu\nu} = \partial_\mu A_\nu - \partial_\nu A_\mu \tag{1.6}$$

so, as required, the longitudinal component of A_μ does not propagate.

So we see that a global symmetry can be made local by adding a new vector particle that mediates an interaction between particles transforming under the symmetry. This is an attractive thing to do because local symmetries are "nicer" than global ones (the associated currents are conserved in all field configurations, unlike global symmetries where the equations of motion must be satisfied), and more importantly, because this seems to correspond to the world we live in.

b) The Standard Model lagrangian

The Standard Model (SM) [15] is based on three local, or gauge, symmetries: the U(1) of hypercharge, weak SU(2) and colour SU(3). These groups act on a variety of fermions, most of which have been discovered. The colour SU(3) only couples to the quarks, which are Dirac fermions (four degrees of freedom) and assumed to come in SU(3) triplets. SU(3) is asymptotically free, since it is a non-abelian gauge theory, so the interaction becomes stronger at large distances and quarks seem to only exist in colour singlet bound states. This makes it difficult, at low energies, to calculate measurable quantities from the lagrangian because it contains quark degrees of freedom and in the real world we see mesons ($q\bar{q}$) and baryons (qqq). However, insofar as it can be tested, quarks and SU(3) colour seem to describe the strong interactions.

The weak SU(2) acts only on the left-handed (LH) fermions (in doublets) so the right-handed (RH) leptons and quarks appear as SU(2) singlets. All fermions carry hypercharge $Y = 2[Q - T_3]$ (Q is electric charge; T_3 is the diagonal SU(2) generator; the photon will therefore be a linear combination of the hypercharge and diagonal SU(2) gauge bosons). The fermions of which matter is composed are

$$\ell_L \equiv \begin{pmatrix} \nu_L \\ e_L \end{pmatrix} , \; e_R \tag{1.7}$$

$$q_L^\alpha \equiv \begin{pmatrix} u_L^\alpha \\ d_L^\alpha \end{pmatrix} , \; u_R^\alpha \; , \; d_R^\alpha \tag{1.8}$$

where α is an SU(3) colour index and runs from 1 to 3. This is the first family, or generation, of the SM. There are two more, consisting of the muon, its associated neutrino and the charm and strange quarks, and then the tau, its neutrino and the top and bottom quarks. The top quark has not yet been seen, but is assumed to exist because of loop effects.

The kinetic term for the fermions is (j is a generation index)

$$\mathcal{L}_f = \sum_j i\left(\bar{\ell}^j_L \not{D} \ell^j_L + \bar{q}^j_L \not{D} q^j_L + \bar{e}^j_R \not{D} e^j_R + \bar{u}^j_R \not{D} u^j_R + \bar{d}^j_R \not{D} d^j_R\right) \tag{1.9}$$

where

$$\not{D} = \gamma^\mu(\partial_\mu - i\frac{\vec{\tau}}{2}\cdot \vec{W}_\mu - i\frac{Y}{2}B_\mu) \tag{1.10}$$

for doublets, and for singlets:

$$\not{D} = \gamma^\mu(\partial_\mu - i\frac{Y}{2}B_\mu). \tag{1.11}$$

$\vec{W}_\mu$ are the triplet of SU(2) gauge bosons, and B_μ is the gauge boson for hypercharge.

The kinetic terms for the gauge bosons are the generalisation of (1.5) and (1.6):

$$\mathcal{L}_{gf} = -\frac{1}{4g'^2}B_{\mu\nu}B^{\mu\nu} - \frac{1}{4g^2}F^A_{\mu\nu}F^{A\mu\nu} \tag{1.12}$$

where g and g' are the hypercharge and SU(2) gauge couplings,

$$B_{\mu\nu} = \partial_\mu B_\nu - \partial_\nu B_\mu \tag{1.13}$$

and $F^A_{\mu\nu}$ is the commutator of the SU(2) covariant derivatives:

$$F^A_{\mu\nu} = \partial_\mu W^A_\nu - \partial_\nu W^A_\mu + \varepsilon^{ABC}W^B_\mu W^C_\nu. \tag{1.14}$$

It is experimentally known that all the fermions (except possibly the neutrinos) have masses, as do the weak gauge bosons. This presents a problem, because mass terms mix the left and right handed components of the fermions, so cannot be gauge invariant in a chiral model. Furthermore, a theory with explicit mass terms for gauge bosons violates unitarity and is not renormalizable because the gauge boson propagator goes to a constant for large momenta, rather than as $1/k^2$.

This problem is solved [16] by realizing that the fermion and gauge boson mass terms are special cases of an interaction with a scalar field (constant field)...and of course Yukawa and scalar kinetic terms can be made gauge invariant. So the trick is to introduce a scalar field with a potential the shape of a mexican hat (ϕ^4 with a negative mass term) so the scalar field will develop a vacuum expectation value (the field sits in the brim of the hat). This destroys the gauge invariance in the Fock space ("Spontaneous Symmetry Breaking") but since the original lagrangian had a gauge symmetry, the theory can be shown to be [17] well behaved and renormalizable.

The weak gauge bosons and the Standard Model fermions get masses from the introduction of a scalar SU(2) doublet called the Higgs (which unfortunately has never been seen)

$$H = \begin{pmatrix} H^+ \\ H^o \end{pmatrix} = \begin{pmatrix} 0 \\ \rho \end{pmatrix} \exp[\frac{i}{2}\vec{\tau}\cdot\vec{\theta}] \tag{1.15}$$

with the hypercharge of the lower component chosen so that it is electrically neutral. (One needs to end up with the unbroken U(1) of electromagnetism, so the scalar that acquires a vacuum expectation value (vev) must be neutral.) Since SU(3) is unbroken, the Higgs carries no colour. The second parametrization of (1.15) will be useful later.

The Higgs lagrangian is

$$\mathcal{L}_H = D_\mu H^\dagger D^\mu H - \frac{\mu^2}{2}H^\dagger H + \frac{\lambda}{4}(H^\dagger H)^2 \tag{1.16}$$

and has a minimum at $< H >= \mu/\sqrt{\lambda} \equiv v/\sqrt{2}$. If the Higgs field configuration is as in equation (1.15) with $\sqrt{2}\rho(x) = v + h(x)$ then one can make a gauge transformation (this cannot affect the physics) to remove the exponent. The physical Higgs is then a real scalar field that corresponds to radial oscillations in the brow of the mexican hat. The three angular degrees of freedom become the longitudinal component of the now massive (the Higgs kinetic term includes $\frac{1}{2}v^2(\vec{\tau}\cdot\vec{W})^2$) gauge bosons.

The Higgs also has gauge invariant Yukawa interactions that give the fermions masses:

$$\mathcal{L}_Y = (h^e_{ij}\bar{\ell}^i H e^j_R + h^u_{ij}\bar{q}^i \hat{H} u^j_R + h^d_{ij}\bar{q}^i H d^j_R) + h.c. \tag{1.17}$$

where h_{ij} is an arbitrary 3×3 matrix of complex coupling constants and $\hat{H} \equiv i\tau_2 H^*$ is the "charge conjugate" Higgs, used so that the up quark can get a mass.

The complete lagrangian for the Standard Model is therefore

$$\mathcal{L}_{SM} = \mathcal{L}_{gf} + \mathcal{L}_f + \mathcal{L}_H + \mathcal{L}_Y \quad . \tag{1.18}$$

1.3 Supersymmetry

a) *Improving the Standard Model: Grand Unified Theories*

The Standard Model has a number of features that are unattractive to theorists [18]: the U(1) charges are not quantized (non-abelian charges are fixed by the commutation relations of the symmetry group algebra), there is no explanation of why there are three generations (there is not even a theoretical definition of generation), there is no relation between the leptons and quarks, the loop corrections to the Higgs mass are quadratic in the cutoff, etc.. Some of these features can be made more to our taste by embedding the SM in a Grand Unified Theory (GUT, see [18]); one assumes that at high energies ($\sim 10^{16}$ GeV) there is a single gauge group (so only one coupling constant) and a complicated scalar sector (there are often more coupling constants here than in the SM), such that the large group (say SU(5)) will spontaneously break to smaller subgroups at some large energy, and this combination of smaller groups will look like the Standard Model when the parameters of the theory are run down from the breaking scale via the renormalization group equations. This solves some of the problems of the SM: the U(1) has quantized charges since it is embedded in a larger group, and the leptons and quarks are partners in the same gauge multiplets in the fundamental theory; they just look very different at low energies because the symmetry is spontaneously broken (like electrons and neutrinos in the SM). Putting the leptons and quarks together in a gauge multiplet has the added attraction that there will be explicit B and/or L violation via the exchange of heavy gauge bosons. This must be very weak at low energies, because the proton has not been observed to decay [19], but could give sufficient B violation in the early Universe to account for the observed excess of matter over anti-matter.

However, there is no explanation of three generations (nothing, I think, can do this) and the loop contributions to the Higgs mass must cancel against the bare mass to $\sim$ 13 decimal places to produce a physical mass $\sim 10^2$ GeV (this si the "hierarchy problem"). Furthermore, there is no relation between the fermions and bosons of the theory; it could be attractive, having unified all the gauge interactions, to represent the bosons and fermions as different realizations of the same field.

There are two possible solutions to the hierarchy problem: put lots of non-perturbative physics above the weak scale (technicolour [20]), or add particles to the theory to cancel the divergences (supersymmetry [11]). Since non-perturbative calculations are difficult, I will concentrate on supersymmetry.

b) More improvements to the Standard Model: supersymmetry

In supersymmetry (SUSY), one introduces a symmetry turning fermions into bosons, which unfortunately doubles the number of particles in the Standard Model because none of the known fermions and bosons can be taken as partners. However, for every divergent scalar loop contribution to the Higgs mass, there is a fermion loop of the opposite sign which will cancel up to terms of order the mass difference between the boson and fermion.

Another reason for looking at supersymmetric theories is that internal, supersymmetric and Poincaré transformations are the only allowed symmetries in field theory. One can show that in a relativistic quantum field theory in four dimensions with a non-trivial S-matrix, the only possible conserved quantities transforming as tensors are the Poincaré generators and the scalar charges of internal symmetries [21]. The (only [22]) loophole that allows one to relate space-time symmetries to supersymmetry in a non-trivial way is that the generators are spin 1/2 Grassman operators, with anti-commutation relations

$$\begin{aligned} \{Q_\alpha, \bar{Q}_{\dot\beta}\} &= 2\sigma^\mu_{\alpha\dot\beta} P_\mu \\ \{Q_\alpha, Q_\beta\} &= \{\bar{Q}_{\dot\alpha}, \bar{Q}_{\dot\beta}\} = \{Q_\alpha, P_\mu\} = \{\bar{Q}_{\dot\alpha}, P_\mu\} = 0 \end{aligned} \tag{1.19}$$

where α and β run from 0 to 1. (Two SUSY transformations give a space-time translation, so local supersymmetry must include gravity!) The $Q_\alpha, \bar{Q}_{\dot\alpha}$ are spin 1/2 Grassman operators, so must take bosons to fermions and vice versa.

One can check [23] that the lagrangian for a free complex scalar and chiral fermion

$$\mathcal{L} = \frac{1}{2}\partial_m A^* \partial^m A + \frac{i}{2}\bar{\psi}\bar{\sigma}^m \partial_m \psi \tag{1.20}$$

is invariant under the infinitesimal SUSY transformation generated by εQ:

$$\begin{aligned} \delta A &= \sqrt{2}\varepsilon\psi \\ \delta\psi &= i\sqrt{2}\sigma^m\bar{\varepsilon}\partial_m A \end{aligned} \tag{1.21}$$

provided that the equations of motion are satisfied. (The fermions in this section are all chiral and in two-component notation—see the appendices of [11] or [18] for conventions, contractions, etc.) This makes sense because an on-shell chiral fermion has two degrees of freedom, but an off-shell one has four (2 complex fields). So one must add another non-propagating scalar field to be the supersymmetric partner of the fermion's off-shell degrees of freedom. The free supersymmetric lagrangian is therefore

$$\mathcal{L} = \frac{1}{2}\partial_m A^*\partial^m A + \frac{i}{2}\bar{\psi}\bar{\sigma}^m\partial_m\psi + \frac{1}{2}F^*F \tag{1.22}$$

and the supersymmetric transformations are

$$\begin{aligned} \delta A &= \sqrt{2}\varepsilon\psi \\ \delta\psi &= i\sqrt{2}\sigma^m\bar{\varepsilon}\partial_m A + \sqrt{2}\varepsilon F \\ \delta F &= i\sqrt{2}\bar{\varepsilon}\bar{\sigma}^m\partial_m\psi \end{aligned} \tag{1.23}$$

From equation (1.22) it is clear that A has mass dimension 1, ψ has mass dimension 3/2, and F dimension 2. ε therefore has dimension -1/2 from equation (1.23), so Q has dimension 1/2 as expected from (1.19).

One can add mass terms and interactions to equation (1.22) without violating the supersymmetry, provided that the superpartners have the same masses and the $\psi\psi A$ vertex has the same coupling constant as AAF. It is this constraint on the masses that forces SUSY, if it exists in the real world, to be broken (we have not seen superpartners for any of the fermions we know).

The less intuitive (but necessary because we use it in chapters 2 and 3) way to do SUSY is to imagine that space-time has four Grassman dimensions ($\theta_\alpha, \bar{\theta}_{\dot{\alpha}}$; $\alpha, \dot{\alpha}$: 1..2) as well as the four real ones we are familiar with. Quantum fields (called superfields) are therefore functions of eight variables, and the $Q_\alpha, \bar{Q}_{\dot{\beta}}$ generate translations in the Grassman dimensions like P_m generates translations in real dimensions. It therefore makes sense that two SUSY transformations are a translation, since in this formalism they are all translations.

Grassman numbers anticommute, so the Taylor expansion of a superfield in θ_α and $\bar{\theta}_{\dot{\alpha}}$ terminates:

$$\begin{aligned} S(x,\theta_\alpha,\bar{\theta}_{\dot{\alpha}}) \quad = \quad & f(x) + \theta\omega(x) + \bar{\theta}\bar{\chi}(x) + \theta\theta m(x) + \bar{\theta}\bar{\theta}n(x) \\ & +\theta\sigma^m\bar{\theta}v_m(x) + \theta\theta\bar{\theta}\bar{\lambda}(x) + \bar{\theta}\bar{\theta}\theta\psi(x) + \theta\theta\bar{\theta}\bar{\theta}D(x) \end{aligned} \tag{1.24}$$

Two things are apparent: there are far more degrees of freedom here than in (1.22), and there is a real vector field in this expansion, so it may be possible to write down supersymmetric gauge theories. The solution to the first problem is to constrain S: if one demands that it be independant of $\bar{\theta}$

$$\frac{\partial}{\partial\bar{\theta}}S(y,\theta,\bar{\theta}) = 0 \tag{1.25}$$

then one has

$$\Phi(y,\theta) = A(y) + \theta\psi(x) + \theta\theta F(y) \tag{1.26}$$

(where y is translated with respect to x of equation 1.24: $y^m = x^m + i\theta\sigma^m\bar{\theta}$ ...not important for handwaving purposes.) Φ is called a chiral superfield, because its fermion component is chiral (LH); it has the right number of degrees of freedom to belong to the lagrangian of equation (1.22). One can also make vector superfields, by requiring $S = S^\dagger$.

For the sake of simplicity, I will forget about vector superfields for the moment, and try to write down a superfield lagrangian that will turn into a sensible real space-time lagrangian for the component fields when we integrate over the Grassman dimensions. Integration is supposed to be a linear operator, so for a Grassman variable η one defines

$$\int \eta d\eta = 1 \tag{1.27}$$

and

$$\int d\eta = 0 \tag{1.28}$$

The only term that will survive when a superfield is integrated over θ and $\bar{\theta}$ is therefore the D-term of equation (1.24). The F-term of a chiral superfield would also survive integration over just θ. Now clearly the product of chiral superfields is a chiral superfield, so if one writes down a "superpotential"

$$P = a_j\Phi_j + m_{jk}\Phi_j\Phi_k + g_{jkl}\Phi_j\Phi_k\Phi_l \tag{1.29}$$

the coefficient of $\theta\theta$ in the Taylor expansion of P will survive the integration over θ. Since the F component field has one higher dimension than the superfield (because $\theta\theta$ has dimension -1), the F-term of P would be dimension four—just right to be a lagrangian. Furthermore, the F-term is a good candidate because we want an action that is invariant under supersymmetry, and δF is a total derivative (see equation (1.23)). One finds as the F-term of P:

$$P\,|_F = a_jF_j + m_{jk}(A_jF_k - \frac{1}{2}\psi_j\psi_k) + g_{jkl}(A_jA_kF_l - \psi_j\psi_kF_l) \tag{1.30}$$

so one can take the $\theta\theta$ component of P and the $\bar{\theta}\bar{\theta}$ component of $P^\dagger$ as an interaction lagrangian. But there are not yet any kinetic terms.

We know that the D-term of a general superfield S has two dimensions more than the superfield. We also know that δD must be a total derivative, because the infinitesimal parameter ε of a supersymmetric transformation has dimension -1/2, so $\delta D \sim \varepsilon(\text{stuff})$ implies that "stuff" has dimension 5/2, and there are no such component fields. One gets kinetic terms for the component fields by taking the D-term of $\Phi\Phi^\dagger$:

$$\Phi\Phi^\dagger\mid_D = i\partial_m\bar{\psi}\bar{\sigma}^m\psi + A^*\Box A + F^*F. \tag{1.31}$$

Solving for F gives

$$\mathcal{L} = P\mid_F + P^\dagger\mid_F + \Phi\Phi^\dagger\mid_D = i\partial_m\bar{\psi}_j\bar{\sigma}^m\psi_j + A_j^*\Box A_j - \mathcal{L}_I \tag{1.32}$$

where

$$\mathcal{L}_I = \frac{1}{2}\sum_{ij}\left(\frac{\partial^2 P}{\partial\phi_i\partial\phi_j}\psi_i\psi_j + h.c.\right) + \sum_i\left|\frac{\partial P}{\partial\phi_i}\right|^2 \tag{1.33}$$

The gauge fields must now be included. By taking suitable derivatives of a vector superfield, one can construct a chiral superfield whose $\theta\theta$ component is the kinetic terms for a massless vector boson (two degrees of freedom) and its associated fermion superpartner (gaugino). The superpotential must be chosen gauge invariant, and the matter field kinetic terms can be made gauge invariant by taking

$$\begin{aligned}\Phi_j^\dagger\exp\{-2gTV\}\Phi_j\Big|_D = \; & F_j^\dagger F_j - D_nA_j^\dagger D^nA_j - i\bar{\psi}_jD_n\bar{\sigma}^n\psi_j \\ & -i\sqrt{2}(A^\dagger T^A\psi\lambda_A - \bar{\lambda}_A\bar{\psi}T^AA) + \tfrac{1}{2}\left|A_j^\dagger TA_j\right|^2\end{aligned} \tag{1.34}$$

where V^A are the vector superfields, $\{T^A\}$ are the generators of the gauge group, D_n is the appropriate covariant derivative, and λ_A are the gauginos. Note that the "kinetic term", or D-term, has introduced a quartic interaction among the scalars, and a Yukawa interaction between a matter scalar and fermion, and a gaugino. This will be useful in chapters 2 and 3.

The supergauge transformation under which equation (1.34) is invariant, and which reduces to an ordinary gauge transformation for the component fields, is

$$\Phi \to \exp\{-iT\Lambda\}\Phi$$

$$\Phi^\dagger \to \Phi^\dagger \exp\{iT\Lambda^\dagger\} \tag{1.35}$$
$$V \to V - \frac{i}{2}(\Lambda + \Lambda^\dagger)$$

where Λ is also a chiral superfield.

1.4 Low temperature B+L violation in the Standard Model

In any gauge theory with SU(2) as a subgroup, there is an infinite set of vacua, labelled by integers that in some sense count how "knotted" the vacuum field configuration is [24]. These vacua are analogous to the countably infinite set of distinct ground states allowed for a particle in a periodic potential if one neglects tunnelling through the barrier. However, in the particle case, one discovers that when tunnelling is included, the correct eigenstates correspond to a particle of fixed momentum ($k \in (0, 2\pi)$) "hopping" along from one minimum to the next. In gauge theory, there are solutions of the Euclidean equations of motion (instantons [25] $\approx$ knots in the gauge field) that allow the field to pass from one vacuum state to another. The true vacua are therefore similar to the momentum eigenstates of the periodic potential, and labelled by an angle θ rather than an integer m.

When fermions are included in the theory, one discovers that gauge field knots act as sources for chiral fermions [26]. Since the SM SU(2) gauge fields couple to LH leptons and quarks, this means that instantons create baryons and leptons (B+L not conserved). Unfortunately tunnelling processes are highly suppressed, so the zero temperature rate could never account for the baryon asymmetry observed in the Universe today [26]. However, at finite temperature, the rate for B+L violation due to these processes is only Boltzmann suppressed ($\sim \exp(-am_W/\alpha_w T)$, $a \simeq 1$) so at temperatures near 100 GeV it exceeds the expansion rate of the Universe and is in thermal equilibrium [2, 3].

a) Instantons and sphalerons

Any finite action field configuration (i.e. something suitable to do perturbation theory about) must go to zero as an appropriate power of $1/r$ at infinity, to within a gauge transformation. Gauge transformations at infinity are mappings from S_3 to the gauge group and can be classified by their winding number [24], which is the number of times the surface at infinity can be wrapped around the gauge group. This is easier to understand for a U(1) theory in two dimensions: the relevant mapping takes the circle at infinity to the angle $\theta \in (0, 2\pi)$ parametrizing the U(1) phase rotation. These mappings can be separated into classes labelled by an integer which is the number of times the mapping goes around the gauge group circle when one makes one circuit around the space-time circle at infinity. It is clear that in this example, mappings with different winding numbers cannot be continuously deformed into each other, so the classes are distinct. This is also the case [24] for an SU(2) gauge group in four space-time dimensions. A vacuum state can therefore be labelled by its winding number, like the position basis states of the periodic potential.

Ordinary particles that satisfy the Schroedinger equation can pass through potential barriers that they classically do not have enough energy to pass over. However, the tunnelling amplitude is small, because the wave function under the barrier is exponentially damped due to the energy being imaginary. Now, one can also get imaginary energy by Wick rotating $t \rightarrow i\tau$, and one can show [24] that by computing the action for solutions of the Euclidean equations of motion that take the particle from one side of the barrier to the other, one can get the tunnelling amplitude.

This is nice, because it can be generalized to gauge field theory. It turns out [25] that there are finite action solutions of the Euclidean equations of motion with winding number ± 1. These (zero-temperature) solutions are called instantons, and correspond to the transformation of an "n" vacuum into an "$n+1$" vacuum. The time-independent lowest energy field configurations (vacua?) are therefore labelled by an angle θ rather than an integer, where θ is analogous to the momentum of the Bloch wave eigenstates of the periodic potential.

However, for $B+L$ violation it is the rate at which the winding number changes that is of interest, not the definition of the vacuum state. In a classical approximation to the path integral, the amplitude to get from an initial to a final state is suppressed by the action for the intermediate field configuration. (One can then compute small corrections about this amplitude as a determinant.) The action for an arbitrary field configuration of winding number $\pm\nu$ has a minimum value, as can be shown via the Cauchy-Schwarz inequality:

$$\begin{aligned} -4g^2 S &= \int d^4x \mathrm{Tr}[FF] \approx \left[\int d^4x \mathrm{Tr} FF \int d^4x \mathrm{Tr} \tilde{F}\tilde{F}\right]^{1/2} \\ &\geq \int d^4x \mathrm{Tr} F\tilde{F} = 32\pi^2|\nu| \end{aligned} \tag{1.36}$$

(Instantons are actually solutions of the euclidean equations of motion that also satisfy $\tilde{F} = F$, so they have action $-8\pi^2/g^2$.) So it is clear that the rate for B+L violation due to field configuration of finite winding number will be infinitesimally small, as it contains a factor $\exp\{-8\pi^2/g^2\}$.

The instanton arguments presented so far have all been at zero temperature. At finite temperature one could hope that new field configurations of unit winding number might appear, corresponding to the field "climbing over the barrier", rather than tunnelling through. This is the case for the particle in the periodic potential, and the finite temperature rate for the particle to climb over is only Boltzmann suppressed ($\sim \exp\{-(\text{height of barrier})/\text{temperature}\}$). It is not obvious that the finite temperature behaviour of large coherent field configurations will mimic that of a particle in a thermal bath [27](far more degrees of freedom in the field theory case). However it has been shown [28] that there are static unstable solutions of the equations of motion that represent the gauge field "sitting on top of the barrier" at T $\neq$ 0. They are called "sphalerons", and the rate per unit volume at which the field climbs over the barrier and passes through the sphaleron configuration can be estimated to be [3]

$$\Gamma \sim \frac{T^4 \omega_s}{m_W(T)} \left(\frac{\alpha_W}{4\pi}\right)^4 N_{tr} N_{rot} \left(\frac{2m_W(T)}{\alpha_W T}\right)^2 \exp\left[-\frac{a(\lambda/g^2) 2m_W(T)}{\alpha_W T}\right] \tag{1.37}$$

where $\omega_s \approx m_W(T)$ is the value of the sphaleron negative energy mode, $N_{tr}N_{rot} \approx 10^3$ comes from integrating over the translational and rotational degrees of freedom and a is of order one. This estimate is only applicable for temperatures below the electroweak phase transition (see [29] for a review of phase transitions in cosmology. The idea is roughly that at finite temperature a particle is slowed down by its interactions with the thermal bath, so it picks up a temperature-dependant contribution to its mass. If the "thermal mass" m_{th} is sufficiently large, the effective mass $\sim -\mu + m_{th}$ is positive and the Higgs vacuum expectation value is zero. The fermions and gauge bosons are therefore massless.); once the gauge bosons become massless, the only dimensionful parameter is the temperature, so one expects a rate [30, 27] per unit volume

$$\Gamma \sim \alpha_W^4 A T^4 \tag{1.38}$$

where $A \sim 10^{-2}$ from numerical simulations [30].

b) The axial anomaly

In a theory of massless fermions, it is classically true that the axial current is conserved:

$$\sum_f \partial_\mu \bar{\psi}_f \gamma^\mu \gamma_5 \psi_f = \partial_\mu J_5^\mu = 0 \tag{1.39}$$

where f is a label that runs over all the fermions in the theory. However, masses violate this conservation law, and a mass scale must necessarily be introduced to renormalize a quantum theory. This leads to the so-called "axial anomaly": although equation (1.39) is classically true, it can be shown that when loop effects are included, one gets

$$\partial_\mu J_5^\mu = \frac{N}{16\pi^2} F\tilde{F} \tag{1.40}$$

Integrating this over a space-time volume between t_1 and t_2 gives

$$Q_5(t_2) - Q_5(t_1) = 2N\nu \tag{1.41}$$

where N the number of fermion generations and ν is the winding number of the gauge field configuration between t_1 and t_2. This is why the rate for instantons or sphalerons is of interest: they emit chiral fermions. In weak SU(2), all the baryons and leptons that couple to the gauge fields are left-handed, so configurations of non-zero winding number must produce a net baryon and lepton number to give $\Delta Q_5 \neq 0$. (They also manage to emit an appropriate combination of fermions to conserve all the gauged currents). So equations (1.37) and (1.38) give the finite temperature rates for $B + L$ violation due to electroweak field configurations of unit winding number. I will frequently refer to these as "sphalerons", although a sphaleron is really an unstable static configuration corresponding to the field sitting on top of the barrier.

c) In the early Universe ...

It is experimentally obvious that at least locally in our corner of the Universe, there is a substantial excess of matter (baryons) over anti-matter. The earth is clearly made of matter, and the sun seems to be, since the solar wind is composed of particles. The galaxy is presumed to be made of matter because we do not see large amounts of γ ray radiation from matter-anti-matter annihilation. This same argument applies to the galaxy clusters we can see, so it appears that if the Universe contains equal amounts of matter and anti-matter, they are separated at least on the scale of galaxy clusters. This is difficult to arrange.

Instead, one can assume that the baryon asymmetry was created at some earlier time. For this to work, one needs a theory [31] with baryon number violation (obviously),C and CP violation (otherwise the same number of baryons and anti-baryons would be generated; B is odd under C and CP), and the interactions must be taking place out of thermal equilibrium, because the equilibrium particle and anti-particle number densities are the same if there are no conserved quantum numbers. It is fairly easy to have interactions out of thermal equilibrium in the early Universe because it is expanding, and the SM contains B and CP violation. However, it was long thought that the anomalous B+L violation of the SM was too weak to be of interest, and that the BAU was generated by the decay of heavy GUT particles. This, of course, could still be true.

The usual definition of thermal equilibrium in the early Universe is that the interaction rate exceeds the expansion rate, or equivalently, that the average time between interactions is less than the age of the Universe. If one uses this criterion for the anomalous $B + L$ violation rates of equations (1.37) and (1.38), one finds that they will be in thermal equilibrium for temperatures satisfying

$$T_c \approx 100 \text{ GeV} < T \leq 10^{12} \text{ GeV}. \tag{1.42}$$

It is not clear whether the BAU can be made via these processes. Numerous models have been proposed [32]. However, it is clear that any B+L asymmetry produced before these interactions come into thermal equilibrium will be wiped out (providing B–L is zero and there are no extra global symmetries; see chapter 4). If, on the other hand, an asymmetry in B–L is created at some temperature $T_B > T_C$ (the GUT scale, for instance), then no interaction violating any combination of B and L independent of that of the "sphalerons" can be in thermal equilibrium between when the asymmetry is created (or when the "sphalerons" come into thermal equilibrium, whichever is later) and when the "sphalerons" drop out of equilibrium (modulo chapter 4). This is the idea of the following chapters, where we will derive constraints on interactions that violate some combination of L and B other than that taken to zero by the "sphalerons".

So for the whole of this thesis, I will assume that the baryon asymmetry was created in some particle physics scenario at a temperature above the electroweak phase transition. Furthermore, this asymmetry must have $B - L \neq 0$ so that it can not be wiped out by the sphalerons alone.

CHAPTER TWO

ESTIMATED CONSTRAINTS

2.1 Introduction

The possibility of generating the cosmological baryon asymmetry via non-equilibrium B-,C-, and CP-violating interactions in the early Universe [31] is a very attractive feature of GUTs. However, an early baryon asymmetry could be washed out by baryon number violating interactions in the later Universe. For instance, non-perturbative electroweak interactions ("sphalerons") [26], which violate $B+L$, are believed to have been in thermal equilibrium in the early universe [2] at temperatures above 100 GeV. The "sphalerons" would have removed any previously existing $B+L$ asymmetry unless $B-L \neq 0$, which is not expected in minimal GUTs such as SU(5). We will assume that the BAU was not made at the electroweak phase transition, although there are a number of models that do this [32].

If there were other perturbative interactions beyond the SM that violated $B-L$, they would wipe out the BAU in conjunction with the "sphalerons". This potential disaster was first explored in the case of $\Delta L = 2$ Majorana neutrino masses by Fukugita and Yanagida [4], who pointed out that it would be averted only if the majorana neutrino masses were < 50 keV. The purpose of this chapter is to extend Fukugita and Yanagida's line of argument to general B and L violating effective interactions, discussing in particular the types of $\Delta B \neq 0$ interactions that could mediate neutron anti-neutron oscillations or $\Delta B \neq \Delta L$ proton decay, and interactions that violate R-parity in supersymmetric theories.

We first consider generic effective B or L violating interations of dimension $2 \leq$

$D \leq 10$, and discuss the requirement that they not take the baryon asymmetry to zero. We will see later (chapter 4) that to be sure to wash out the BAU, one must have all the Standard Model interactions in equilibrium with the "sphalerons" and the extra B or L violating operator. Some of the Higgs couplings are quite weak, so the lightest right-handed fermions do not come into chemical equilibrium until fairly low temperatures, and any 'beyond the Standard Model' interactions could be in equilibrium until this happens if there is an asymmetry stored in the RH fermions. We therefore list, in tables A.1, A.2, A.3 and A.4, bounds calculated using a variety of maximum temperatures on the standard catalogue of SU(3) × SU(2) × U(1) invariant B and/or L violating operators made out of fields present in the Standard Model [34, 33] and its supersymmetric extension [35, 33]. Similar results can be derived in a model including a singlet neutrino (super)-field. For each operator, we tabulate the best previous limit on its magnitude as well as the new limits we obtain from the persistence of the cosmological baryon asymmetry. We then work out the applications of this general analysis to $n - \bar{n}$ oscillations, R-parity breaking, and the decay of the lightest supersymmetric particle (LSP). Our limits suggest that $\tau_{n-\bar{n}} \gg 10^8$ sec (using $T_{max} \sim T_{GUT}$), and $\tau_{LSP} \gg 10^{-10}$ sec in most models, so that it seems very unlikely that neutron-anti-neutron oscillations or LSP decay will be observed in present laboratory experiments. It remains possible that the LSP could decay on a timescale relevant to cosmology.

2.2 Constraints from cosmology

Let us first consider the general case of an effective non-renormalizable interaction of dimension $D = 4 + n$

$$\mathcal{L}_D = \frac{\mathcal{O}_D}{M^n} \tag{2.1}$$

where M is a mass parameter (but not a mass scale, because there are phase space factors and coupling constants mixed into it) and $n \gtrsim 1$. We assume that the generation structure of the interaction is generic, and does not accidentally preserve one of the e, μ or τ lepton flavours. Quark generations are not conserved in the SM and hence cannot provide an accidental conserved quantity that might vitiate our argument [5]. The corresponding interaction rate is

$$\Gamma_D(T) \sim \frac{T^{2n+1}}{M^{2n}} \tag{2.2}$$

for $T \leq M$, and thermodynamic equilibrium is avoided if $\Gamma_D(T) \leq H(T)$, where H is the Hubble expansion rate:

$$H \sim 25T^2/m_{pl} \tag{2.3}$$

giving

$$M^{2n} \gtrsim \frac{m_{pl}T^{2n-1}}{25} \tag{2.4}$$

In the case of a $\Delta B = 2$ interaction that could cause $n - \bar{n}$ oscillations, persistence of the BAU requires that the out-of-equilibrium condition (2.4) be imposed for all temperatures below that at which the baryon asymmetry is generated (probably, see [36]). This would be $\sim 10^{14}$ GeV in the conventional GUT scenario, in which case we would require

$$M > 10^{14+2/n} \text{ GeV} \tag{2.5}$$

for $\Delta B = 2$ interactions.

Let us assume the existence of strong non-perturbative B and L violating but $B-L$ conserving electroweak interactions with an interaction rate $\Gamma_{sph} \sim 0.01\alpha^4 T\,\theta(T-T_c)$ [30], where $T_c \sim 100$ GeV is an effective threshold temperature whose precise value depends on details of the electroweak theory (see section 4c of the introduction).

Since the BAU could have been generated at any temperature between the electroweak phase transition and m_{pl}, and since the SM interactions that ensure the washout of the baryon asymmetry drop out of equilibrium at various temperatures, we will consider a series of conditions for the preservation of the BAU: absence of thermodynamic equilibrium for any 'beyond-the Standard-Model' $B - L$ violating interaction at temperatures (A) T = T_c = 100 GeV, (B) T= 10^5 GeV, (C) T= 10^7 GeV, (D) T= 10^9 GeV, and (E) T= 10^{12} GeV. We present rough estimates of the constraints here. In the next chapter we will calculate the interaction rates more carefully and list the bounds in the tables at the end. These more accurate bounds are slightly weaker than those we present here.

As one might expect for dimension $D > 4$ operators, the bounds are significantly better at higher temperatures. We find from equation (2.4) that

$$M > T_{max} \left(\frac{10^{18}}{T_{max}} \right)^{\frac{1}{2n}} \quad . \tag{2.6}$$

The parametrization (2.2) is not applicable to renormalizable interactions. We estimate the rate for an operator of the form

$$\mathcal{L}_D = \lambda O_4 \tag{2.7}$$

to be $\Gamma_4 \sim \lambda^2 T$. This is out of equilibrium if

$$\lambda^2 < \left(\frac{25}{m_{pl}} \right) T \tag{2.8}$$

so that the most stringent upper bound on λ comes from case(A): T $\sim$ 100 GeV

$$\lambda \lesssim 10^{-8} \tag{2.9}$$

We parametrize superrenormalizable D = 3 and D = 2 interactions by coefficients m and μ^2 respectively, yielding interaction rates

$$\Gamma_3 \sim \frac{m^2}{T}, \quad \Gamma_2 \sim \frac{\mu^4}{T^3} \tag{2.10}$$

For these two types of interaction, the most stringent upper bounds on the parameters m and μ again come at $T \sim 100$ GeV:

$$m \leq 10^{-6}\ \text{GeV}, \qquad \mu \leq 10^{-2}\ \text{GeV} \tag{2.11}$$

The constraints (2.6), (2.9) and (2.11) have dramatic implications for models, as we now discuss.

2.3 Standard Model B and L violating operators

The effective operators describing new interactions beyond the SM must be made out of light fields, and invariant under the symmetries which survive below the scale M, including Lorentz invariance, the appropriate gauge and discrete symmetries, and perhaps supersymmetry. Often these residual symmetries suffice, by themselves, to rule out the appearance of operators of a given dimension which violate some global conserved charge. For example, it is impossible to construct gauge and Lorentz invariant renormalizable B or L violating interactions within the Standard Model. This is no longer true in the supersymmetric version of the theory, where the addition of squarks and sleptons permits [35] the construction of renormalizable superpotential terms that violate B or L. In most models, these are removed by the imposition of an additional discrete symmetry, such as R–parity [38]. If one removes the requirement of supersymmetry, as in the soft SUSY-breaking terms induced via supergravity [39], then further D = 2 or 3 B and L violating operators are possible.

In table A.3, we display the types of operators that may appear at various dimensions involving the Standard Model fields. The notation is schematic: q, l represent quarks or leptons (q^c and l^c are their conjugates), h is the Higgs scalar, and $\mathcal{D}$ is a gauge-covariant derivative. All indices are suppressed. There are various ways of producing gauge and Lorentz singlets ... for example, $qqql$ represents six independent operators formed from different combinations of chirality, weak isospin, and hypercharge.

The $D = 5$ entry in the table is the effective low energy lepton number violating operator arising from the see-saw neutrino mass generation mechanism [40], whose cosmological implications were first discussed by Fukugita and Yanagida. The remaining operators involve baryon number violation, and laboratory limits on these operators involve an estimate of their hadronic matrix elements. We use the estimates of Costa and Zwirner [33] for nucleon decay matrix elements to convert nucleon decay partial width limits into estimates of the operator mass scale listed in the column under "laboratory limit".

For the $D = 7$ operators which could induce $\Delta B = -\Delta L$ proton decay, we find that the cosmological limits from the preservation of the BAU at $T_{max} = 10^{12}$ GeV yield limits on M that are stronger than direct laboratory searches by four orders of magnitude. This is even more true for B and L violating operators of dimension $D=$ 9, 10, or 11.

The $D = 9$ six-quark operator violates baryon number by two units without violating L. So the laboratory limit arises from the bound on neutron-anti-neutron oscillations [41] ($\tau_{n \to \bar{n}} > 1.2 \times 10^8$ seconds) which we convert to a limit on M using matrix element estimates from the reviews [33] and [42]. Cosmologically, the violation of baryon number alone means that the effects of this operator may (or may not [36]) be able to relax the baryon number of the Universe to zero in the absence of the "sphalerons". This depends on the combination of left and right quark fields, as they carry different fractions of the BAU. It will be discussed in chapter 4. Assuming the BAU does go to zero, the limits in the columns of our table should be interpreted as the lower bound on M (a mass scale divided by coupling constants) necessary to ensure the survival of baryon number generated at $T \sim T_{max}$.

2.4 Supersymmetric B or L violating operators

We now take a general look at sources of baryon and/or lepton number violation in supersymmetric theories. With the common definition of R-parity: $R = (-1)^{2S+L+3B}$, violation of B and/or L usually also signifies the breaking of R–parity (S is the spin). R-parity violation may appear directly in exactly supersymmetric theories, e.g. those with additional superpotential terms or D-terms beyond the minimal extension of the Standard Model.

The MSSM superpotential is

$$P_{MIN} = h_u H Q U^c + h_d \bar{H} Q D^c + h_e \bar{H} L E^c + \varepsilon H \bar{H} \qquad (2.12)$$

where the chiral multiplets are denoted by H and $\bar{H}$ for the Higgs doublets, Q and L for the left-handed quark and lepton doublets, and U^c, D^c and E^c for the right-handed singlet fermions. The coupling h_i are the standard Higgs Yukawa couplings and ε is a Higgs mass mixing parameter necessary to avoid a massless axion-like state. R-parity could be introduced explicitely by adding to the superpotential

$$\begin{aligned} P_R = \; & m_1 LH + \lambda_1 LLE^c + \lambda_2 LQD^c + \lambda_3 U^c D^c D^c \\ & + \frac{1}{M_1} QQQL + \frac{1}{M_2} QQQ\bar{H} + \frac{1}{M_3} U^c U^c D^c E^c \\ & + \frac{1}{M_4} HH\bar{H}L + \frac{1}{M_5} HHLL + \frac{1}{M_6^2} U^c D^c D^c U^c D^c D^c + \ldots \end{aligned} \qquad (2.13)$$

Alternatively, sources of R-parity breaking may be found in D-terms or soft SUSY breaking terms, which take the form

$$\begin{aligned}\mathcal{L} \ni \; & \mu_1^2 \tilde{l} h + \mu_2^2 \tilde{l} \tilde{h}^\dagger \\ & + m_2 \tilde{l} \tilde{l} \tilde{e}^c + m_3 \tilde{l} \tilde{q} \tilde{d}^c + m_4 \tilde{u}^c \tilde{d}^c \tilde{d}^c \\ & + m_5 l \tilde{h} + m_6 \bar{\tilde{h}}^\dagger\end{aligned} \tag{2.14}$$

The explicit R-parity-breaking terms we consider are listed in tables A.1, A.2 ($D \leq 4$) and A.4 ($D > 4$). In addition the generation of a sneutrino ($\tilde{\nu}$) vev could break R-parity spontaneously [44]. As in the case of the non-supersymmetric Standard Model described above, the survival of the cosmological baryon asymmetry in spite of such terms will impose much tighter constraints than are available from laboratory cross-section measurements and the absence of rare processes.

As was discussed in detail by Hall and Suzuki [45], the $m_1 LH$ term is a source for Majorana neutrino masses, and the strongest constraint they derive:

$$m_{1e} \leq 0.1 \frac{\varepsilon}{30\ \text{GeV}}\ \text{GeV} \tag{2.15}$$

is based on the mass limit ($m_{\nu_e} \leq 1$ eV) from neutrino–less double beta decay (We have rescaled their limit to account for the more recent constraint on the Majorana neutrino mass.) From equation (2.11) we see that our (admittedly rough) limit is five orders of magnitude stronger. It is also generation independent, whereas the bound from neutrino masses is much weaker for ν_μ and ν_τ. Analogous arguments place constraints on the soft masses: $\mu_{1e} \leq O(1)$ GeV from Majorana neutrino masses and $\mu_1 \leq 10^{-2}$ GeV from our arguments. The lab bound is again weakened by going to higher generations.

The dimension-4 terms in (2.12) have been discussed extensively in the literature [43, 46, 56, 47, 48, 49, 50]. The strongest constraints on λ_1 and λ_2 are on the off-diagonal elements [50, 43, 46, 56]:

$$\lambda_1^{113} \leq 2 \times 10^{-3} \left(\frac{\tilde{m}}{100\text{GeV}}\right)^{1/2} \tag{2.16}$$

$$\lambda_2^{113} \leq 10^{-3} \left(\frac{\tilde{m}}{100\text{GeV}}\right)^{1/2} \tag{2.17}$$

due to their radiative contribution to the electron neutrino Majorana mass. Charged current universality of quark and lepton couplings to W bosons leads to the limits

$$\lambda_1^{12k} \leq 0.04 \left(\frac{\tilde{m}}{100\text{GeV}}\right) \tag{2.18}$$

$$\lambda_2^{12k} \leq 0.03 \left(\frac{\tilde{m}}{100\text{GeV}}\right) \tag{2.19}$$

and the absence of neutrino-less double beta decay gives

$$\lambda_2^{111} \leq 3 \times 10^{-3} \left(\frac{\tilde{m}}{100\text{GeV}}\right)^{5/2} \tag{2.20}$$

The λ_3^{11k} coupling is strongly constrained by the absence of $n - \bar{n}$ oscillations as was first pointed out by Zwirner [47], who found

$$\lambda_3^{11k} \leq 2 \times 10^{-7} \left(\frac{\tilde{m}}{100\text{GeV}}\right)^{5/2} \tag{2.21}$$

(see also reference [49] where a stronger limit is derived from $NN \to KK + \ldots$, but is subject to larger nuclear uncertainties). The remaining bounds in [56] are typically O(.1) to O(.4) from a variety of sources such as charged π decay, τ decay, $\nu_\mu - e$ scattering, etc. If two or more such couplings are non-zero, typical constraints [48, 43] $\lambda < 10^{-3}$ may be obtained from $\mu \to e\gamma$ or $\mu \to 3e$. From equation (2.9), we find the surprisingly strong cosmological estimate

$$\lambda_{1,2,3} \lessdot 10^{-8} \tag{2.22}$$

which is again generation independent.

We are not aware of any bounds on the soft terms $m_{2,3,4}$. These are generic $D = 3$ mass terms, so we find

$$m_{2,3,4} \leq 10^{-6}\ \mathrm{GeV} \tag{2.23}$$

We cannot improve existing bounds on the $D = 5$ supersymmetric F-term $QQQL$ because it involves the same fields as the effective operator induced by the "sphalerons". We are not aware of any previous bound on $M_2^{-1}QQQ\bar{H}$, but we find here using $T_{max} = 10^5$ GeV that $M_2 > 10^{11}$ GeV. This is of course analogous to the bound on $M_5^{-1}HHLL$ already derived by Fukugita and Yanagida. Finally for the $D = 7$ term, the current limit from $n - \bar{n}$ oscillations is

$$M_6 > 10^8 \left(\frac{100\mathrm{GeV}}{\tilde{m}}\right)^{2/3}\ \mathrm{GeV} \tag{2.24}$$

whereas we find

$$M_6 > 10^{4.7} \left(\frac{T_{max}}{100\mathrm{GeV}}\right)^{5/6}\ \mathrm{GeV} \tag{2.25}$$

Lastly, we come to limits on $\tilde{\nu}$ vevs which would break R-parity spontaneously [44]. In the absence of any explicit R-parity breaking terms in the superpotential, spontaneous R-parity breaking would lead to a massless goldstone boson, the majoron. For majoron emission from stars to be compatible with observation, one must have [51] a sneutrino vev of less than 9 keV. Explicit R-parity breaking would remove the astrophysical bound, but not our cosmological constraints which require the vev to be orders of magnitude below the weak scale (assuming of course, that the vev forms above T = 100 GeV).

A sneutrino vev would induce mixing between the corresponding neutrino and gauginos leading to a non-vanishing neutrino mass [45]. Upper limits on the $\tilde{\nu}$ vev [48] from Majorana neutrino masses range from O(1) MeV for $< \tilde{\nu}_e > \neq 0$ to O(10) GeV for $< \tilde{\nu}_\tau > \neq 0$.

A sneutrino v.e.v. in the ordinary Yukawa term $h_e\bar{H}LE^c$ leads to a $D = 3$ lepton number violating term $h_e < \tilde{\nu} > \tilde{h}e^c$. However, we get better constraints from the D-term $g < \tilde{\nu} > \lambda\ell_L$, where λ is the gaugino. We find that for all generations

$$m \leq 3 \times 10^{-6} GeV \tag{2.26}$$

We do not address here the technical question of whether such a small v.e.v. can be generated in a realistic model before the non-perturbative electroweak interactions drop out of equilibrium. This is of course a necessary condition for our constraints, and it is what we would expect if the sneutrino v.e.v. were comparable in size to that of the Higgs. Such a large v.e.v. is ruled out by our bounds. The maximum size of a sneutrino v.e.v. that does not cohabit with the "sphalerons" is a dynamical question that we do not address here.

2.5 What can we see in the lab?

We comment finally on the observability of neutron-anti-neutron oscillations and R-parity breaking in the light of our results. In the case of the D = 9 $qqqqqq$ operator, to tolerate $M \sim 10^6$ GeV close to the laboratory limit we require the baryon asymmetry to be created at a temperature below

$$T \sim \left(\frac{25}{m_{pl}} M^{2n}\right)^{\frac{1}{2n-1}} \simeq 10^5 \text{ GeV} \tag{2.27}$$

This suggests that if $n - \bar{n}$ oscillations are seen, then the baryon asymmetry was made at the EPT (or maybe produced as an e_R asymmetry; see chapter 4). A similiar calculation in the MSSM for the D = 7 $U^cD^cD^cU^cD^cD^c$ operator gives $T \leq 10^6$ GeV. This would be consistent with a RH lepton asymmetry produced at high temperatures and later turned into a baryon asymmetry by the "sphalerons", with baryogenesis at the EPT and with the Affleck and Dine [79] scenario.

If R-parity was broken, then the lightest supersymmetric particle (LSP) could decay. However, this would be observable in the laboratory only if the LSP decay path length were ≤ 1 m, corresponding to a life-time $\tau_{LSP} \leq 10^{-8}$ seconds. There are two important candidate decay modes for the LSP: LSP $\rightarrow \nu \bar{f} f$ and LSP $\rightarrow \nu\gamma$. One can easily estimate that the former has a partial life-time [49, 43]

$$\tau_{\nu f \bar{f}} \simeq \frac{10^{-21}\ \text{sec}}{\lambda^2} \left(\frac{100\text{GeV}}{m_{LSP}}\right)^5 \left(\frac{m_{\tilde{f}}}{100\ \text{GeV}}\right)^4 \tag{2.28}$$

where λ is one of the trilinear superpotential couplings in equation (2.12). An analogous estimate gives

$$\tau_{\nu\gamma} \simeq \frac{10^{-17}\ \text{sec}}{\lambda^2} \left(\frac{100\text{GeV}}{m_{LSP}}\right)^3 \left(\frac{m_{\tilde{f}}}{100\ \text{GeV}}\right)^4 \tag{2.29}$$

We found that $\lambda \leq 10^{-8}$, which from equations (2.28) and (2.29) suggests that the decay of the LSP would be unobservable.

Similiar estimates apply, as one can verify after making a basis change, if one breaks R-parity by an HL coupling or a sneutrino vev. This suggests that R-parity breaking will be unobservable in laboratory experiments, modulo loopholes discussed in more detail in the following chapters.

We conclude that the survival of the cosmological baryon asymmetry imposes severe constraints on $B-L$ violating interactions beyond the Standard Model. These bounds are in many cases far stronger than those derived from laboratory experiments. A corollary is that many interesting proposed non-Standard-Model phenomena (R-parity violation,$n \leftrightarrow \bar{n}$ oscillations) appear to have little chance of being seen in the laboratory in the near future if our assumptions are correct.

CHAPTER THREE

MORE DEPENDABLE CONSTRAINTS

3.1 Introduction

It appears that non-perturbative electroweak interactions that violate baryon number B and lepton number L, whilst conserving the combination $B - L$, are likely to have been strong in the early Universe at temperatures above about 100 GeV [2],[3]. They therefore pose a threat to the baryon asymmetry of the Universe, if it was generated primordially at higher temperatures, as in many scenarios for baryogenesis. This threat is not fatal if the early Universe had asymmetries in $B - L$ and $B + L$, since these could not both be wiped out by the non-perturbative electroweak interactions. However, in this case the addition to the Standard Model of any interactions violating some other combination of B and L would, if they were in thermal equilibrium in the early Universe, be fatal for the survival of any primordial baryon asymmetry.

This important observation was first made by Fukugita and Yanagida[4], who used it to constrain Majorana neutrino masses (see also Refs. [5],[6] which extend the Fukugita-Yanagida analysis). Their argument has subsequently been used on many other conjectured extensions of the Standard Model [7, 8],[9], including interactions that violate R parity in supersymmetric extensions of the Standard Model, and $\Delta B = 2$ interactions that could lead to neutron-antineutron oscillations or nucleon-nucleon annihilations in nuclei. (A related study of $\Delta B = 2$ limits had also appeared in Ref. [43].)

R parity could be violated either explicitly by B and/or L-violating interactions in the lagrangian[45], and/or spontaneously by a vacuum expectation value (v.e.v.) for a sneutrino field: $< 0|\tilde{\nu}|0 > \neq 0$[44]. The cosmological bounds would apply to any source of explicit R-violation, or to spontaneous R-violation if it were present at some temperature T above $T_c \sim 100$ GeV, below which the rates for non-perturbative electroweak interactions are presumably negligible[2]. *A priori*, these constraints on explicit R-violation appeared so severe as to make it seem very unlikely that a violation of R-parity could be seen in an accelerator experiment[7, 8]. Likewise, the constraints on putative $\Delta B = 2$ interactions may (this depends on how the BAU is made) make it difficult for $n - \bar{n}$ oscillations or $N - N$ annihilations in nuclei to be observable in the laboratory[7, 8].

However, we should emphasize that there are several escape routes from these discouraging arguments. The baryon asymmetry might be generated at some low temperature $T < T_c$. R parity might be violated spontaneously by a v.e.v. that only appears at some $T < T_c$. The B- and/or L-violating interactions might conserve[5],[7] some other quantum number able to "encode" the primordial baryon asymmetry and hence allow it to survive strong non-perturbative effects. (This "conserved" quantum number could be associated with an exact symmetry of the lagrangian [1], or could be "effectively conserved" by the interactions in chemical equilibrium, but violated by some other out-of-equilibrium operator.)

The purpose of this chapter is to develop more thouroughly our cosmological constraints on B and L violation, and to confirm the claims of the previous chapter. In Section 2 we consider carefully the rates for interactions due to couplings of various different dimensions, compare them with the Hubble expansion rate, and infer the corresponding cosmological upper limits on the couplings. After including the numerical factors due to thermal averaging, etc., we find bounds which are far stronger

[1] This is possible in the leptonic sector, to the extent that the neutrinos do not mix, but not in the quark sector, where there is substantial Cabibbo-Kobayashi-Maskawa mixing between the different generations.

than many laboratory constraints. In Section 3 we revisit constraints on explicit and spontaneous lepton number violation[4]-[9]. In Section 4 we address R parity, first giving a complete discussion of possible R-violating effects, including the possibilities of spontaneous R violation and a right-handed neutrino. Then we give more precise bounds on the coefficients of these interactions, incorporating the results of Section 2. These are then translated into lower limits on the lifetime of the lightest supersymmetric particle (LSP). We find that the LSP lives long enough to escape from a laboratory experiment before decaying, and may well live long enough to be subject to astrophysical constraints on particles decaying after nucleosynthesis[52],[53]. In this case it is likely to live longer than the age of the Universe[54], which would enforce very stringent upper bounds on possible R-violating interactions. R-violating decays of the Z^0[55], such as $Z^0 \rightarrow \tau +$ (chargino) should be unobservable, as should other signatures of R-violation at future hadron-hadron colliders such as the LHC or SSC[46]-[57]. The only exception to these arguments is a small window for the LSP to be charged or strongly-interacting, and sufficiently short-lived to escape the constraints from cosmological nucleosynthesis, in which case the LHC, SSC or LEP might see massive, apparently stable particles instead of the usual missing energy signature for supersymmetry. At the end of this section we also discuss briefly cosmological aspects of an alternative[58] to R parity in supersymmetric models. In Section 5 we update our previous discussion[7] of $\Delta B = 2$ interactions.

3.2 Rate estimates and bounds on couplings

a) Interaction rates

In this section we evaluate the reaction rates, taking into account the thermal averages appropriate for early Universe calculations, and then estimate the bounds on the different types of couplings.

We begin by considering the simplest case of a $D = 4$ interaction with dimensionless coupling, which could be either of Yukawa or quartic scalar type. In a supersymmetric theory, which is our main interest, these would both arise from a superpotential term of the generic form

$$F \ni h_{ijk}\Phi_i\Phi_j\Phi_k \tag{3.1}$$

for some chiral superfields $\Phi_{i,j,k}$. The quartic scalar couplings arise from

$$|\frac{\partial F}{\partial \Phi_i}|^2 \ni h_{ijk}^2|\phi_j\phi_k|^2 \tag{3.2}$$

(we denote by ϕ the spin-0 components of the chiral superfields Φ: the fermionic components will be denoted by ψ) which yields the simple squared matrix element

$$|\mathcal{M}|^2 = h^4 \quad . \tag{3.3}$$

However, such quartic couplings do not violate any quantum numbers, since the same particles are destroyed then re-created. The one-to-two decay processes arising from the same superpotential term (3.1):

$$\frac{\partial^2 F}{\partial \Phi_j \partial \Phi_k}\psi_j\psi_k \ni h_{ijk}\phi_i\psi_j\psi_k \tag{3.4}$$

do violate quantum numbers and yield a squared matrix element

$$|\mathcal{M}|^2 = 2h^2 p\cdot p' \tag{3.5}$$

for fermions with final-state four-momenta p, p'. The thermally-averaged decay rate is then given by

$$\Gamma_{12} = \int \frac{d^3p_0}{(2\pi)^3 2E_0}\cdot\int\frac{d^3p}{(2\pi)^3 2E}\cdot\int\frac{d^3p'}{(2\pi)^3 2E'}(2\pi)^4\delta^{(4)}(p_0 - p - p')|\mathcal{M}|^2\frac{f_0}{n_0} \tag{3.6}$$

where the subscript 0 refer to the initial-state spin-0 particle with thermal distribution

$$f_0 = \left[e^{E_0/T} - 1\right]^{-1} \tag{3.7}$$

and number density

$$n_0 = \frac{\zeta(3)}{\pi^2}T^3 \ . \tag{3.8}$$

Evaluation of the integrals in (3.6) yields

$$\Gamma_{12} = \frac{h^2 m_0^2 \pi}{192\zeta(3)T} = 1.4 \times 10^{-2}\frac{m_0^2 h^2}{T} \tag{3.9}$$

at high temperatures $T \gg m_0$, the decaying scalar mass.

In the case of $D = 3$ interactions, we can consider superpotential terms of the form

$$F \ni m_3^{ij}\Phi_i\Phi_j \tag{3.10}$$

or soft supersymmetry-breaking lagrangian terms of the forms

$$\mathcal{L} \ni -m_3^{S^{ijk}}\phi_i\phi_j\phi_k - m_3^{F^{ij}}\psi_i\psi_j \ . \tag{3.11}$$

We will consider the latter two possibilities (3.11) here. In the case of the three-scalar process of scalar decay, the high-temperature rate is similar to that in (3.9), namely

$$\Gamma_{12}^0 = \frac{m_3^2\pi}{192\zeta(3)T} \simeq 0.014\ m_3^2/T \ . \tag{3.12}$$

The rate due to the supersymmetry-breaking fermion mass term in (3.11) is best treated by first rediagonalizing the fermion mass matrix. For off-diagonal terms $\Delta m_f^{ij} \simeq m_3^{F^{ij}} \ll m_f$, the resulting basis rotation will be by an angle $\theta \simeq \Delta m_f/m_f \simeq m_3^F/m_f$. In the rotated basis there are species-changing effective gauge vertices that alter the external species through the mixing angle θ. The resulting interactions have effective rates given by θ^2 times the normal gauge interaction rates, and hence of order

$$\Gamma_{12}^{\Delta} \simeq \theta^2\Gamma_{12} \simeq \frac{\pi g^2\theta^2 T}{192\zeta(3)} \simeq 1.4 \times 10^{-2}\ g^2\frac{(m_3^F)^2}{(m_f)^2}\ T \tag{3.13}$$

which we will typically use at $T \sim m_f$, so that it becomes of order $1.4\times10^{-2}g^2(m_3^F)^2/T$, where g may represent any of the Standard Model gauge couplings.

Finally, $D = 2$ interactions could also arise from soft supersymmetry-breaking terms in the Lagrangian:

$$\mathcal{L} \ni -\mu_2^{2^{ij}}\phi_i\phi_j \tag{3.14}$$

which can be regarded as mass insertions on external scalar lines. Again, if we change basis to diagonalize the squared-mass matrix, we get an effective interaction which has the form of a standard gauge interaction, which changes species by an amount of order $\theta \sim \Delta m^2/m_0^2 \simeq \mu_2^{2^{ij}}/m_0^2$. The interaction rate is therefore of order

$$\Gamma_{12}^{\Delta 0} \simeq \theta^2\Gamma_{12}^0 \simeq 1.4\times10^2\frac{(\mu_2^2)^2}{m_0^4}\,g^2T \tag{3.15}$$

which we will typically use at $T \sim m_0$, so that it becomes of order $1.4\times10^{-2}(\mu_2^2)^2g^2/T^3$.

Finally, we consider non-renormalizable interactions O_D of dimension $D = 4+n$, whose strengths we parametrize by mass parameters M_D:

$$\mathcal{L}_D = \frac{O_D}{I_D^n}\ . \tag{3.16}$$

There are large numbers of possible interactions, and for each one the numbers of initial- and final-state particles can be varied (e.g., for an interaction of six fields, 1 $\rightarrow$ 5, 2 $\rightarrow$ 4, 3 $\rightarrow$ 3, etc.). We do not attempt to perform accurately the appropriate higher-order thermal averages, but simply estimate

$$\Gamma_D = \frac{10^{-3}}{M_D^{2n}}\cdot T^{2n+1} \tag{3.17}$$

where the numerical coefficient of Γ_D has been chosen to agree roughly with the previous calculations (3.9),(3.12).

b) Comparison with Hubble expansion rate

We now compare the interaction rates with the Hubble expansion rate:

$$H = \sqrt{\frac{8\pi}{3}\rho G_N} \tag{3.18}$$

where we approximate the energy density by a free gas of massless particles:

$$\rho \simeq \frac{\pi^2}{30} N_{eff} T^4 \tag{3.19}$$

and N_{eff} is the effective number of degrees of freedom for the minimal supersymmetric extension of the Standard Model

$$N_{eff} = (N_B + \frac{7}{8} N_F) = \frac{915}{4} \quad . \tag{3.20}$$

Substituting the number (3.20) into (3.19) and (3.18), we find

$$H \simeq 25 \frac{T^2}{m_{pl}} \tag{3.21}$$

to be compared with the reaction rates Γ of the previous section.

We take as our criterion for maintaining the cosmological baryon asymmetry the simple out-of-equilibrium condition $\Gamma(T) < H(T)$ for

$$T_c < T < T_{max} \leq T_M \simeq 10^{12} \text{ GeV} \tag{3.22}$$

where, in the absence of any detailed model, we assume that the non-perturbative electroweak interactions drop out of equilibrium at $T_c = 100$ GeV. T_M is an estimate of the temperature at which they may have *started* to be in equilibrium, assuming that they had a rate[30]

$$\Gamma_{sph} \sim (0.01 \text{ to } 0.1)\alpha^4 T \tag{3.23}$$

Taking $T_{max} = T_M$ is likely to be too stringent a constraint for two reasons. The first is that "sphalerons" only take $(B + L)_L$ to zero. Since the RH fermions will be out of chemical equilibrium down to $T_R \ll T_M$ (see the next chapter), it could be possible to store the asymmetry in some right-handed particle and have various $B - L$ violating operators in equilibrium above T_R. Secondly, it is well known that cosmological baryosynthesis would have been possible with reduced efficiency even if the analogous Sakharov[31] out-of-equilibrium condition for B-, C- and CP-violating interactions were not respected. For comparison, we note here that the efficiency of baryosynthesis would exceed 10^{-3} if the ratio of the B violating interaction rate to the expansion rate is of order 10[60]. Aggressive model-builders may feel free to exploit a similar grace factor here, if they need it.

The out-of-equilibrium condition (3.22) becomes, for the different interactions discussed in Section 3.2 a),

$$25\frac{T^2}{m_{pl}} > \left\{ \begin{array}{l} 1.4 \times 10^{-2}\frac{m_0^2 h^2}{T},\ 1.4 \times 10^{-2}\frac{m_3^2}{T},\ 1.4 \times 10^{-2} g^2 (m_3^F)^2/T, \\ \qquad 1.4 \times 10^{-2} g^2 (\mu_2)^4/T^3,\ \frac{T^{2n+1}10^{-3}}{M^{2n}} \end{array} \right\} \tag{3.24}$$

from Eqs. (3.9), (3.12), (3.13), (3.15), (3.17) respectively. It is evident that the strongest restrictions on the $D \lesssim 4$ interactions come from $T \sim T_c \simeq 100$ GeV. For orientation, we take the unknown mass-scale $m_0 \sim T_c \simeq 100$ GeV also, although this could be an underestimate by an order of magnitude in which case our bounds on h from (3.9) and μ_2^2 from (3.15) should also be relaxed by a factor of three and an order of magnitude, respectively. In the case of (3.13), we take $m_f = m_{\chi^\pm}$, $g = g_2$ for orientation. In the cases of the non-renormalizable $D > 4$ interactions (3.17), the strongest restrictions come from the maximum temperature we can impose bounds at. This is the minimum of T_M (see equation 3.22), the temperature at which the BAU is created, and, if the asymmetry is stored in a RH fermion, the temperature at which it goes out of equilibrium (see chapter 4). Tables A.1 and A.2 compile all the upper limits on renormalizable B- and/or L-violating interactions that we extract from Eq. (3.24). Tables A.3 and A.4 list the bounds we can compute for higher dimensional operators at various temperatures. Also shown are the laboratory upper limits on each of these couplings[43],[45],[46],[56],[47]-[50]. We see that in many cases the cosmological bounds (3.24) improve on the laboratory bounds by many orders of magnitude. This has the consequence that many of the laboratory searches for physics beyond the Standard Model are doomed to failure if our cosmological argument applies. However, we should re-emphasize that there are loopholes in this argument, as discussed in paragraph 4 of Section 3.1.

3.3 Application to lepton number violation

a) Explicit lepton number violation and neutrino masses

The above argument that non-perturbative electroweak phenomena, perturbative B- and/or L-violation from some source beyond the Standard Model, and the survival of a primordial cosmological baryon asymmetry were mutually incompatible was first developed by Fukugita and Yanagida[4]. They applied this argument to a $\Delta L = 2$ Majorana mass term for neutrinos, using it to derive bounds on neutrino masses[4]-[6]. In this section we first recapitulate the essence of their argument to exemplify the quantitative discussion of the previous section, and then go on to consider models with spontaneous violation of lepton number, and hence Majorons[61].

A $\Delta L = 2$ Majorana mass term

$$\mathcal{L}_M = -m_M \nu\nu \tag{3.25}$$

is of the $D = 3$ type (3.11) introduced earlier, and hence bounded by the third constraint in (3.24), as already noted in the Table:

$$m_M < 3 \times 10^{-5} \text{ GeV} \ . \tag{3.26}$$

It is not entirely clear that this bound is directly applicable to a majorana neutrino mass, because it was derived for a particle with a large L conserving mass and a small L violating mass correction. To conservatively avoid the confusion arising from having L violating propagation but L conserving interactions, one can compute the rates for interactions with an L violating internal propagator. The rate for the decay $W^- \to e^-\nu$ (ordinary W decay with to $e\bar{\nu}$ with a majorana mass insertion on the neutrino leg) is

$$\Gamma \simeq 3 \times 10^{-3} \frac{g^2 m_M^2}{T} \tag{3.27}$$

for small m_M. This gives a similar bound to (3.26).

A large right-handed ν_R Majorana mass term M_M is also of interest. Indeed, the requirement that Majorana neutrino masses and electroweak non-perturbative effects do not erase a pre-existing baryon asymmetry, is the complement of a previous insight of Fukugita and Yanagida[62], that the *out*-of-equilibrium decay of a ν_R respecting such conditions could generate a lepton number asymmetry, which non-perturbative electroweak effects could convert into the observed baryon asymmetry.

It is clear that the L violating ν_R propagation would be out of equilibrium if

$$M_M \gtrsim T_M \simeq 10^{12}\ \mathrm{GeV} \quad . \tag{3.28}$$

However, equation (3.28) may again not be the relevant bound. The particles involved in the L violating interaction (in this case, ν_R propagation) are not those in the effective $B+L$ violating vertex induced by "sphalerons", so the $B-L$ asymmetry among this second group of particles will survive if they are not in thermal equilibrium with the ν_R. As with the left-handed majorana neutrino masses, it is probably safer to compute an L violating interaction rate among Standard Model particles induced by the heavy ν_R. In the limit $M_M \gg T$, the rate for $e_i^+ H^- \to e_j^- H^+$ is

$$\Gamma_{ij} \approx 3 \times 10^{-4} \left|\frac{h_{ik} h_{jk}^*}{M_k^2}\right|^2 T^3 \quad . \tag{3.29}$$

Requiring $\Gamma_{ii} < H$ gives bounds on the sum of masses divided by coupling constants:

$$\sum_k \frac{h_{ik} h_{ik}^*}{M_k} < \sqrt{\frac{10^5}{m_{pl} T_{max}}} \tag{3.30}$$

where T_{max} is the maximum temperature at which we want these interactions out of equilibrium. Fortunately, this is the same combination of masses and coupling constants that enters the "LH" neutrino seesaw mass matrix. If we assume that we are working in a basis where this is diagonal, we have that the left-handed majorana neutrino masses induced by the seesaw satisfy:

$$m_M < \sqrt{\frac{10^5}{m_{pl} T_{max}}} \frac{<H>^2}{2} \simeq 100 \text{ keV} \sqrt{\frac{1 \text{ TeV}}{T_{max}}} \quad . \tag{3.31}$$

Using $T_{max} \sim 1$ TeV, we therefore infer that

$$m_\nu \lesssim \max (50 \text{ keV}, \ 100 \text{ keV}) \quad . \tag{3.32}$$

This argument does not exclude a left-handed Majorana origin for the possible 17 keV neutrino mass claimed recently[63], and is of course completely compatible with the neutrino mass difference of order 10^{-3} eV that is the most appealing explanation[64] for the apparent solar neutrino deficit[65], tolerating either a left- or right-handed Majorana origin for neutrino masses of this order.

b) Spontaneous lepton number violation and Majorons

There has been much interest in the possibility that lepton number might be an exact global symmetry which is broken spontaneously by a Higgs field with $L = 2$, in which case a Goldstone boson emerges, the Majoron[61]. Models with $I = 1$ triplet Majoron fields now appear to have been ruled out by LEP data on the invisible Z^0 decay width[66], leaving only the $I = 0$ singlet Majoron possibility. In the MSSM, the scalar partner of the neutrino (which has $L = 1$) could also acquire a vev. The constraints on the associated majoron are considerably stronger than those on the singlet because the sneutrino is a doublet; this will be discussed later under R_P violation. If one introduces a singlet scalar ϕ with couplings to singlet neutrinos N^i (RH) via the lagrangian

$$\partial_\mu \phi \partial^\mu \phi + \lambda_i \phi N^i N^i + h.c. \tag{3.33}$$

and ϕ acquires a vev

$$\phi(x) = [V + \eta(x)] \exp i\{\frac{\chi}{V}\} \tag{3.34}$$

then the first term of the expansion of the exponential in the Yukawa interaction gives a Majorana mass $M_i = \lambda_i V$ to the N_i, and the second term is

$$i\lambda_j \chi (N^j N^j - \bar{N}^j \bar{N}^j) \ . \tag{3.35}$$

If one adds a Dirac mass matrix between the SM neutrinos ν_L^k and the N^j with $m_D \ll \lambda V$, then diagonalizing the seesaw mass matrix mixes a small amount of N into ν_L:

$$\nu \simeq \nu_L + \frac{m_D}{M} N \tag{3.36}$$

so that the effective coupling of the singlet majoron χ to neutrinos is of the form

$$\mathcal{L}_\chi = \frac{i}{2} \sum_{j,k} \bar{g}_{jk} \chi \nu_j \gamma_5 \nu_k \tag{3.37}$$

where the entries in the matrix g_{jk} are of order

$$\bar{g}_{ij} \sim \frac{m_j \text{ or } m_k}{V} \ . \tag{3.38}$$

One might guess that V is at least as large as the conventional $I = \frac{1}{2}$ Higgs v.e.v.: $V \gtrsim$ 100 GeV. The strongest previous upper bound on the g_{jk} comes from consideration of the supernova 1987a. As argued in Ref.[67], the remnant neutron star core is probably filled for $t_\nu \sim 10$ seconds with a degenerate ν_e gas. In this case, one can estimate the energy emission rate per unit time and volume due to $\bar{\nu}_e\nu_e \rightarrow \chi\chi$ annihilation to be

$$q \simeq \frac{p_F^5}{1152\pi^5} \sum_j g_{ej}^4 (3 \ln \frac{4p_F^2}{m_{\nu_e}} - \frac{23}{2}) \tag{3.39}$$

resulting in the upper bound

$$\bar{g}_e \equiv (\sum_j g_{ej}^4)^{\frac{1}{4}} \lesssim 1.5 \times 10^{-6} \frac{(E_\chi/2 \times 10^{53} erg)^{\frac{1}{4}} \, (R/10km)^{\frac{1}{2}}}{(t_\nu/10s)^{\frac{1}{4}} (N_{\nu_e}/2 \times 10^{57})^{\frac{5}{12}}} \tag{3.40}$$

where E_χ is the neutron star binding energy emitted in the form of Majorons, R is the radius of the neutron star, and N_{ν_e} is the number of ν_e's in the core.

The cosmological argument of Fukugita and Yanagida [4] can be used to strengthen this bound considerably. Since we expect $V \gtrsim 100$ GeV, which is supported by the astrophysical bound (3.40), we could expect that the large singlet v.e.v. V forms at some temperature T_0 before the electroweak phase transition, i.e., above T_c. In this case, there is an epoch $T_c < T < T_0$ during which apparent L violation and non-perturbative electroweak B- and L-violation coexist. The L violating rate for $\chi\nu \rightarrow \bar{e}W$ is approximately

$$\Gamma_M \sim 10^{-4} |\bar{g}_{jk}|^2 T \tag{3.41}$$

which is slightly smaller than the decay rate for $W \rightarrow \chi\nu\nu$:

$$\Gamma_W \sim 10^{-3} \frac{g^2 g_{jk}^2 m_W^2}{T} \ . \tag{3.42}$$

Requiring these interactions to be less than the Hubble expansion rate gives

$$\bar{g}_{jk} \lesssim 10^{-6} \tag{3.43}$$

Similar bounds apply to the analogous Majoron couplings of the μ and τ generations.

3.4 Application to R-violation

a) Classification of R-violating interactions

As we have already indicated, the presence of non-perturbative electroweak interactions that violate B and L but conserve $B-L$ severely constrains other interactions that violate $B-L$. In view of the standard definition of R parity[68] in supersymmetric theories:

$$R = (-1)^{2S+L+3B} \tag{3.44}$$

where S is the spin, sources of R violation must also violate either L and/or B.

R parity may be broken either explicitly[45] and/or spontaneously[44], and may or may not involve fields beyond those in the Minimal Supersymmetric Standard Model. The latter can be described by the superpotential

$$F_{SM} = h_u HQU^c + h_d \bar{H}QD^c + h_e \bar{H}LE^c + \epsilon H\bar{H} \tag{3.45}$$

where the chiral supermultiplets are denoted by H and $\bar{H}$ for the Higgs doublets, $Q(L)$ for the left-handed quark (lepton) doublets, and U^c, D^c and E^c for the right-handed singlets in the model. The couplings h_i are the standard Higgs Yukawa couplings, and ϵ is a Higgs mixing parameter. Note that we have suppressed all generation indices, and all the matter supermultiplets (couplings h_i) should be understood as vectors (matrices) in generation space. R-violation could be introduced explicitly[45] by adding to the superpotential

$$\begin{aligned} F_{RX} &= m_1 LH + \lambda_1 LLE^c + \lambda_2 LQD^c + \lambda_3 U^cD^cD^c \\ &+ \frac{1}{M_1}QQQL + \frac{1}{M_2}QQQ\bar{H} + \frac{1}{M_3}U^cU^cD^cE^c \\ &+ \frac{1}{M_4}HH\bar{H}L + \frac{1}{M_5}HHLL + \frac{1}{M_6^3}U^cD^cD^cU^cD^cD^c + \ldots \end{aligned} \tag{3.46}$$

Alternatively, sources of R-violation may be found in soft supersymmetry-breaking terms, which take the form

$$\mathcal{L} \ni \mu_1^2 \tilde{\ell} h + \mu_2^2 \tilde{\ell} \tilde{\bar{h}}^+ + m_2 \tilde{\ell} \tilde{\ell} \tilde{e}^c + m_3 \tilde{\ell} \tilde{q} \tilde{d}^c + m_4 \tilde{u}^c \tilde{d}^c \tilde{d}^c + m_5 \ell \tilde{h} + m_c \ell \tilde{\bar{h}}^+ \tag{3.47}$$

where the lower-case letters (with tildes) denote the Standard Model particle (spartner) members of the supermultiplets in (3.45). We do not consider any special generation structure for the matrix couplings $(m_i, \lambda_i, \mu_i, M_j)$, but are aware that our subsequent arguments could in principle be evaded by invoking specific generation-dependent symmetries [5], [7], [36], [37].

As already mentioned, the alternative to explicit R-breaking is spontaneous R-violation due to a v.e.v. for one of the sneutrinos $\tilde{\nu}$ in one of the L supermultiplets[44]. A $\tilde{\nu}$ v.e.v. may develop with or without any of the above explicit R-violating terms[69]. Our cosmological bounds would apply to spontaneous R-violation that persisted above the electroweak phase transition. Within the context of the minimal supersymmetric Standard Model (3.45) without explicit R-violation, the sneutrino vev would give rise to an fermion mass term of the form

$$F_{SM} \ni h_e \ < \tilde{\nu} > \tilde{h} e_R^c \tag{3.48}$$

which would be subject to cosmological bounds of the same type as those applying to the explicit R-violating term $m_1 LH$ in (3.46). Another source of effective L-violating vertices in the presence of $< \tilde{\nu} > \neq 0$ would be the $SU(2)$ gaugino interaction:

$$\mathcal{L} \ni \sqrt{2} g_2 < \tilde{\nu} > \tilde{W}^+ \ell^- + h.c. \tag{3.49}$$

which is an effective $D = 3$ fermion mass-mixing term of the type shown in (3.11).

In the absence of explicit R-violation, the spontaneous breakdown of R-parity would lead to a massless Goldstone boson, the Majoron. This Majoron is the phase of an $SU(2)$ doublet, so is severely constrained by various laboratory and astrophysical observations. In particular, since the sneutrino is a doublet, its associated majoron mixes slightly ($\sim<\tilde{\nu}>/<H>$) with the Higgs, and thereby acquires a direct coupling to ordinary matter. Requiring this coupling to be small gives the (severe) upper bound $<\tilde{\nu}>\lesssim 9$ keV [51]. Explicit R-violating terms are often added to avoid this constraint.

Another possible source of explicit and spontaneous (via $\tilde{\nu}_R$ v.e.v.) R violation [69] is to add a right-handed neutrino singlet field ν_R via the following terms in the superpotential:

$$F_{\nu_R} = m_N \nu_R \nu_R + h_N H L \nu_R + \lambda \nu_R^3 \quad . \tag{3.50}$$

The above are the simplest mechanisms for R-violation; the most relevant laboratory bounds[43],[45],[46],[56] on these sources of R-violation are compiled in the Tables. Also shown are the constraints on these R-violating couplings that arise from our cosmological bounds (3.24). We see that the cosmological bounds are considerably stronger, at least for renormalizable interactions. This has dramatic consequences for the search for R-violating signatures of supersymmetry that we discuss in the next section, subject of course to the caveats that we mentioned in paragraph 4 of Section 3.1.

b) Lightest supersymmetric particle decay

When sources of R-parity violation are introduced, the lightest R-odd particle (the Lightest Supersymmetric Particle or LSP) becomes unstable. In most supersymmetric models the LSP is a neutral spin-$\frac{1}{2}$ fermion, called a neutralino, which is some linear combination of the neutral gauginos and higgsinos[70]. To compute the decay rate and branching ratios of the LSP, we consider the case of explicit R violation, and work in the mass eigenstate basis in which the possible R-violating fermion mass term from an HL superpotential term has been rotated away[45], and all the R-violating fermion interactions now arise from the following trilinear superpotential terms:

$$^3F_{RX} = \lambda_1 LLE^c + \lambda_2 LQD^c + \lambda_3 U^c D^c D^c \tag{3.51}$$

where we have suppressed the implicit generation indices. We note that the trilinear R-violating superpotential terms do not involve the Higgs superfields.

The most important mechanisms for LSP decay are associated with tree and one-loop diagrams. We consider first tree-level LSP decay. Since neither higgsinos nor gauginos participate in the R-violating superpotential (3.44), LSP decay proceeds through a first coupling to a fermion-sfermion pair, followed by virtual sfermion decay via an R-violating trilinear interaction. Since the higgsino-fermion-sfermion coupling is related to a small Higgs-matter Yukawa coupling, tree-level LSP decay normally proceeds mainly from the gaugino components of the neutralino. This might not be true if the higgsino components of the LSP dominate over the gaugino components, or if the LSP decay to a top quark is kinematically accessible and the R-violating $U^cD^cD^c$ superpotential term is non-negligible for the top, or if the ratio $\tan\beta = v_2/v_1$ of supersymmetric Higgs v.e.v.'s is very large and the b quark Yukawa coupling is large. However, we shall ignore these possibilities. We will also ignore final-state particle masses, and therefore treat fermions of different helicities independently. In the limit of degenerate sfermion masses, we find for the LSP$\rightarrow f_1 f_2 f_3$ decay width

$$\begin{aligned} \Gamma \;=\; & \frac{\lambda^2 M_{LSP}}{256\pi^3}\Bigg\{4(g_1^2+g_2^2+g_3^2)\Bigg[\frac{1}{16}-\frac{3(M_{LSP}^2-m_{\tilde f}^2)}{8M_{LSP}^2} \\ & +\frac{(3m_{\tilde f}^2-M_{LSP}^2)(M_{LSP}^2-m_{\tilde f}^2)}{8M_{LSP}^4}\ln\left(\frac{m_{\tilde f}^2}{m_{\tilde f}^2-M_{LSP}^2}\right)\Bigg] \\ & +(g_1g_2+g_1g_3+g_2g_3)\,\frac{m_{\tilde f}^2}{M_{LSP}^2}\Bigg[\ln\left(\frac{m_{\tilde f}^2}{m_{\tilde f}^2-M_{LSP}^2}\right)\left(\frac{m_{\tilde f}^2}{M_{LSP}^2}+\right. \end{aligned} \tag{3.52}$$

$$\frac{m_{\tilde{f}}^2}{M_{LSP}^2}\ln\left(\frac{m_{\tilde{f}}^2}{2m_{\tilde{f}}^2 - M_{LSP}^2}\right) - 1\Big) - 1 + \frac{m_{\tilde{f}}^2}{M_{LSP}^2}\left(\sum_{n+1}^{\infty}\frac{m_{\tilde{f}}^{2n} - (m_{\tilde{f}}^2 - M_{LSP}^2)^n}{(2m_{\tilde{f}}^2 - M_{LSP}^2)^n n^2}\right)\Big]\Big\}$$

$$\simeq \frac{\lambda^2}{6144\pi^3}\frac{M_{LSP}^5}{m_{\tilde{f}}^4}\left[g_1^2 + g_2^2 + g_3^2 + g_1 g_2 + g_2 g_3 + g_3 g_1\right] \quad (m_{\tilde{f}} \gg m_\chi) \tag{3.53}$$

where λ is the relevant R-violating superpotential term from Eq. (3.51), and g_i is the coupling of the i'th fermion-sfermion supermultiplet to the combination of gaugino components in the LSP. There is an identical rate for decay to the CP-conjugate channel, and to obtain the total LSP decay rate one must sum over all the different matter supermultiplets in all generations that can participate in the LSP decay via R-violating interactions.

Violation of R parity may also induce LSP decay to two-body final states via loop diagrams. Colour conservation forbids decays to quarks and gluons. Decays to $H^+\ell^-$ or $H^0\nu$ could only proceed via two-loop diagrams. The LSP can decay to $W^\pm\ell^\mp$ or $Z\nu$ if it is heavy enough, and the rates for these decay modes are likely to be comparable to the $\gamma\nu$ decay mode whose rate we now compute.

Since we are only interested in decay into an on-shell photon, only the dipole piece of the induced electromagnetic vertex for LSP $\rightarrow \gamma\nu$ contributes. Parametrizing the induced dipole vertex as

$$^D\mathcal{M}_\alpha = -i\bar{u}(p_\nu)\,\frac{i\sigma_{\alpha\beta}q^\beta}{m_{LSP} + m_\nu}\,(F_2^V + F_2^A\gamma_5)u(p_{LSP}) \tag{3.54}$$

the LSP decay rate is[71]

$$\Gamma(LSP \rightarrow \nu + \gamma) = \frac{m_{LSP}}{8\pi}\,(1 - \frac{m_\nu}{m_{LSP}})^2\,(1 - \frac{m_\nu^2}{m_{LSP}^2})\left[|F_2^V|^2 + |F_2^A|^2\right] . \tag{3.55}$$

Computation of the loop diagrams gives, for an LSP that is dominated by its photino component:

$$F_2^V = F_2^A = -\frac{e^2\lambda}{(4\pi)^2\sqrt{2}}\,M_{LSP}(I_1 + I_2) \tag{3.56}$$

where I_1 and I_2 are parametric integrals whose leading terms in an expansion in $m_{LSP}/m_{\tilde{f}}$ are

$$I_1 \simeq -\frac{1}{6}\frac{m_{LSP}}{m_{\tilde{f}}}, \quad I_2 \simeq -\frac{1}{2}\frac{m_{LSP}}{m_{\tilde{f}}} \quad . \tag{3.57}$$

Plugging the approximations (3.57) into the expressions (3.56) for the dipole vertex factors F_2^{V}, and then into the decay rate formula (3.55), we find

$$\Gamma_{Loop} \simeq \frac{\alpha^2\lambda^2}{512\pi^3}\frac{m_{LSP}^5}{m_{\tilde{f}}^4} \tag{3.58}$$

where λ is the relevant R-violating superpotential coupling from Eq. (3.51). Again, there is also an equal rate for decay to the CP-conjugate channel, and one must sum over all the possible final-state neutrino generations and over the flavours of matter supermultiplets that can appear in the loop.

To compare the competing decay rates (3.58) and (3.53), we make the simplifying assumption that one type of R-violating coupling dominates, and we assume also that the gaugino couplings in (3.53) have the magnitude

$$(g_1^2 + g_2^2 + g_3^2 + g_1 g_2 + g_2 g_3 + g_3 g_1) \simeq 4\pi \cdot \frac{\alpha}{\sin^2\theta_W} \tag{3.59}$$

so that

$$\frac{\Gamma_{tree}}{\Gamma_{Loop}} \simeq \frac{\pi}{3\sin^2\theta_W\alpha} \simeq 600 \quad . \tag{3.60}$$

We therefore conclude that the tree-level three-body decays are likely to dominate in LSP decay.

To get some idea of what the LSP lifetime might be, we use equation (3.59) as before, obtaining

$$\tau_{LSP} \simeq \left(\frac{10^{-6}}{\lambda}\right)^2 \left(\frac{20 \text{ GeV}}{m_{LSP}}\right)^5 \left(\frac{m_{\tilde{f}}}{200 \text{ GeV}}\right)^4 \times 10^{-4} s \tag{3.61}$$

where we have scaled the important unknown quantities by our cosmological upper limit on λ, and values of m_{LSP} and $m_{\tilde{f}}$ that are closer to experimental lower limits in the minimal supersymmetric extension of the Standard Model. The estimate (3.61) will be compared with astrophysical[52],[53] and accelerator constraints in the next subsections.

c) Astrophysical limits on LSP decays

We base our analysis on the updated astrophysical constraints on massive unstable relic particles given in Ref. [53]. There upper bounds were given on $m_{LSP}(n_{LSP}/n_\gamma)$ as a function of the LSP lifetime. The standard calculations[70],[72] of the relic LSP abundance are still largely applicable in the presence of R-violating interactions respecting the cosmological constraints discussed in Section 2 of this chapter. Typically, they give n_{LSP} comparable to or even larger than n_B, so that stable LSP's could provide the critical density for closure of the Universe, $\Omega_{LSP} \equiv \rho_{LSP}/\rho_c = 1 : \rho_c \sim 2 \times 10^{-29} h_0^2 gcm^{-3}$ where the present Hubble expansion rate $H_0 \equiv h_0 \times 100 kms^{-1} M_{pc}^{-1}$. We focus for the purposes of orientation on the likelihood that unstable LSP's have the density comparable to n_{LSP}^c which would have given $\Omega_{LSP} = 1$ if they had been stable. We consider illustrative cases where $n_{LSP} = n_{LSP}^c$ $(100 n_{LSP}^c)(0.01 n_{LSP}^c)$. Examination of Figs. 3 and 6 of Ref. [53] then indicates that the LSP lifetime should either obey

$$\tau_{LSP} \lesssim 1s \ (0.1s) \ (1000s) \tag{3.62}$$

or

$$\tau_{LSP} \gtrsim 10^{17} y \ (10^{19} y) \ (10^{15} y) \quad . \tag{3.63}$$

The upper bound (3.62) comes from considerations of the effects of LSP decay hadron showers on the primordial abundances of light nuclei, and the lower bound (3.63) from upward-going decay muons in underground detectors. These constraints would be respected if

$$\lambda \gtrsim \left(\frac{20\ \text{GeV}}{m_{LSP}}\right)^{5/2} \left(\frac{m_{\tilde{f}}}{200\ \text{GeV}}\right)^2 \times 10^{-8}(10^{-7.5})\ (10^{-9.5}) \tag{3.64}$$

or if

$$\lambda \lesssim \left(\frac{20\ \text{GeV}}{m_{LSP}}\right)^{5/2} \left(\frac{m_{\tilde{f}}}{200\ \text{GeV}}\right)^2 \times 10^{-20}(10^{-20.5})\ (10^{-19.5})\ . \tag{3.65}$$

We see that there is roughly three orders of magnitude of λ values allowed between the cosmological upper bound $\lambda \lesssim 10^{-6}$ (see the Table) and the astrophysical lower bound (3.64), whose width depends on the masses of the LSP and the sfermions as well as on n_{LSP}. Assuming for definiteness that $n_{LSP} = n^c_{LSP}$, we see that the window is open only for

$$m_{LSP} \gtrsim 20\ \text{GeV} \left(\frac{m_{\tilde{f}}}{2\ \text{TeV}}\right)^{4/5} \tag{3.66}$$

Using the model-independent LEP lower bound $m_{\tilde{f}} \gtrsim 45$ GeV[73], we infer that the window near $\lambda \sim 10^{-6}$ is open only for

$$m_{LSP} \gtrsim 1\ \text{GeV}\ \ . \tag{3.67}$$

If Nature does not exploit this window, she must respect the stringent upper bound (3.65), which corresponds to $\lambda \lesssim (\frac{m_W}{m_P})$, and is very suggestive that R parity is in fact strictly conserved.

d) Comparison with Z^0 decay limits

If R parity is not conserved, supersymmetric particles may be produced singly in either hadron-hadron collisions or e^+e^- annihilation. In view of the cleanliness and large number of Z^0 decays at LEP, these provide the most stringent accelerator upper limits on R-violation available for the time being. The Z^0 could decay into an ordinary fermion and a supersymmetric one (a gaugino/higgsino mixture), provided the sum of their masses is less than m_{Z^0}. In the case of the LSP, the lower bound (3.61) on τ_{LSP} indicates that it would decay outside any detector. Hence, if the LSP is neutral as generally expected, its production in association with a neutrino would simply add to the invisible decay width of the Z^0. This would be detectable only if the mixing between the neutralino and a neutrino were relatively large. Using the latest experimental determination[66] of the number of equivalent neutrinos in Z^0 decay, $N_\nu = 2.99 \pm 0.05$, one can only infer the following upper bound on the LSP/ν mixing angle $\theta_{LSP/\nu}$:

$$\sin\theta_{LSP/\nu} < 0.2 \quad (95\ \%\ c.l.) \tag{3.68}$$

if one assumes for simplicity that the neutralino components of the LSP have negligible couplings to the Z^0, which is strictly true for any linear combination of higgsinos.

The Z^0 could also decay to a SM charged lepton and a supersymmetric fermion.The cleanliness of this reaction means that a branching ratio of 10^{-7} could be detectable when LEP gets 10^7 Z^0's, corresponding to a mixing angle below 10^{-3}. In an interaction-eigenstate basis, the Z^0 necessarily decays into a particle-antiparticle pair. However, since the propagating fermions must be mass eigenstates, it is possible for the Z^0 to decay into an apparently ordinary charged lepton (a τ, for instance) and a (mostly) supersymmetric chargino (which we denote by χ) if there are R-violating contributions to the charged fermion mass matrix[55]. In a one-generation model, the superpotential terms that could give such contributions are

$$F \ni m_1 HL + h_\tau \bar{H} L E^c \tag{3.69}$$

with a non-vanishing sneutrino v.e.v. $< 0|\tilde{\nu}|0 > \equiv \tilde{V}$ in the second term. Including D-term contributions and an $SU(2)$ gaugino mass M_2 gives a charged fermion mass matrix

$$(\tilde{W}^-, \tilde{h}_1^-, \tau^-) \begin{pmatrix} M_2 & \sqrt{2} m_W \cos\beta & 0 \\ \sqrt{2}\, m_W \sin\beta & -\epsilon & h_\tau \tilde{V} \\ g\tilde{V} & -m_1 & m_\tau \end{pmatrix} \begin{pmatrix} \tilde{W}^+ \\ \tilde{h}_2^+ \\ \tau^+ \end{pmatrix} \qquad (3.70)$$

where $\tan\beta = v_2/v_1$.

The matrix (3.70), which we denote by M, can be diagonalized by independent unitary transformations U_L and U_R:

$$U_L^+ M U_R = M_{diag} = diag(M_+, M_-, m_\tau) \ . \qquad (3.71)$$

Since sparticles have not been pair-produced at LEP, we know that $M_\pm > 45$ GeV[73]. In fact, in the area of the $\epsilon - M_2$ plane where $M_- > 45$ GeV, one finds $M_+ > 90$ GeV, so the heavier sparticle can be ignored. Denoting the relevant left-eigenvectors of M by ℓ_-, ℓ_τ, the relevant right-eigenvectors by r_-, r_τ, and defining the matrix

$$A \equiv \frac{g}{\cos\theta_W} \left[T_3 - \sin^2\theta_W Q_{em}\right] \qquad (3.72)$$

the effective coupling of the Z^0 to the $\tau^+\chi^-$ mass eigenstates will be

$$|A_{eff}|^2 = |\ell_\tau^+ A \ell_-|^2 + |r_\tau^+ A r_-|^2 \qquad (3.73)$$

yielding a branching ratio for $Z \to \tau^+\chi^-$ of

$$B(Z^0 \to \tau^+\chi^-) = 0.72(1 - \frac{M_-^2}{m_Z^2})^2 \, (1 - \frac{M_-^2}{3m_Z^2})|A_{eff}|^2 \qquad (3.74)$$

to which should be added the corresponding branching ratio for decays into the CP-conjugate state.

What if LEP should see $Z^0 \to \tau\chi$ events with a sample of, say, 10^7 Z^0 decays? We show in the figures contours of the values of m_1 and $\tilde{V}$ that would give $B(Z^0 \to \tau^+\chi^-) = 5 \times 10^{-7}$ and hence 10 $Z^0 \to \tau\chi$ events. Three parameter choices are shown: Fig. B.1 is for $m_1 \neq 0$, $\tilde{V} = 0$, corresponding to explicit, not spontaneous R-violation, Fig. B.2 is for $m_1 = 0$ and $\tilde{V} \neq 0$, corresponding to spontaneous, not explicit R-violation, and Fig. B.3 is for a mixed case.

We see that the typical values of m_1 required in the figures to get an observable rate for $Z^0 \to \tau\chi$ decay are many orders of magnitude larger than those allowed by the cosmological upper bound of Section 2. Thus, if LEP should see any $Z \to \tau\chi$ events, one would conclude that our cosmological argument was being evaded, *either* because the baryon asymmetry of the Universe was generated after the electroweak phase transition, *or* because R-violation occurs spontaneously after the electroweak phase transition, *or* because some accidental symmetry allows the baryon asymmetry to survive non-perturbative effects. In the absence of such deviousness, our cosmological argument would have led us to expect mixing $\theta_{\chi/\tau}$ between the supersymmetric fermion and the τ of order

$$\theta_{\chi/\tau} \sim \frac{g\tilde{V} \text{ or } m_1}{m_\chi - m_\tau} \lesssim 10^{-8} \tag{3.75}$$

and hence $B(Z^0 \to \tau\chi) \lesssim 10^{-16}$.

e) Implications for LHC physics

The possible signatures of R-violation at high-energy *pp* colliders have been discussed in Ref. [46]-[57]. As discussed there, three classes of R-violating events can be distinguished:

1) R conservation in sparticle pair-production, R violation in their decays,

2) R violation in single sparticle production, R conservation in decay,

3) R violation in single sparticle production, R violation in decay.

Modulo the caveats expressed in the last paragraph of the previous section, the cosmological upper bounds on R-violation derived in Section 2 suggest the following prognoses for observing these classes of events.

1) In this case, the sparticle pair-production cross-section would be as large as in normal analyses. The lower limit (3.61) indicates that any R-violating decays could only occur outside the detector. Hence, if the LSP is neutral as usually expected, the signal for supersymmetry is the same missing-energy signature normally expected for R-conserving models[68],[70]. However, if the magnitude of R-violation is in the window (3.64), the normal cosmological argument[70] that the LSP *must* be weakly-interacting and neutral (based on the non-observation[74] of anomalous heavy nuclei) is no longer valid, and one could entertain the possibility that the LSP is strongly-interacting or charged, even though this occurs rarely in models. A strongly-interacting LSP would be confined with conventional quarks and/or gluons inside neutral or charged hadrons. A heavy neutral hadron would be difficult to distinguish from a conventional neutron or K_L^0 in a detector. A heavy, apparently stable charged particle, either strongly- or only electromagnetically-interacting, would be easy to see in principle[2]. We conclude that sparticles could be seen via this class of events, although possibly via the unusual signature of massive "stable" charged particles.

2) In this case, the bounds of Section 2.1 suggest that R-violating production cross-sections will be too small to be detected. By the time the LHC or SSC comes into operation, LEP will have provided model-independent lower limits on sparticle masses of order m_W or m_Z. Hence the cross-sections for single sparticle ($\tilde{X}$) production will

[2]Such a particle is not excluded by unsuccessful searches for anomalous heavy isotopes, because these are only sensitive to stable relics with lifetimes $> 10^{10}y$.

be suppressed relative to single W or Z production by a factor

$$\frac{\sigma(\tilde{X})}{\sigma(W)} \lesssim \frac{\lambda^2}{g_2^2}\left(\frac{3}{2J_{\tilde{X}}+1}\right) \lesssim 10^{-12} \tag{3.76}$$

Since the LHC or SSC is expected to produce at most $0(10^9)W$ particles, we conclude (modulo the caveats at the end of the previous section on Z^0 decay) that this class of events should be unobservably rare, although the missing energy or "stable" charged particle signature would be detectable in principle. Conversely, observation of such a class of events would indicate that Nature was evading our cosmological arguments by one of the mechanisms mentioned earlier.

3) According to our cosmological arguments, the production cross-section for this class of events should be as unobservably small as for the previous class, and there is the added disadvantage that it is likely to be more difficult to pick out an R-violating decay signature such as $\tilde{X} \to L+$ jet or jet + jet.

We conclude that, if our cosmological arguments are correct, the prospects for observing R violation at the LHC are not bright. The only, somewhat distant, possibility of a distinctive signature appear to be that the LSP lives in the window (3.64) and is strongly-interacting or charged, offering the possibility of detecting heavy, apparently stable charged sparticles.

f) Constraints on generalized matter parity

It has recently been pointed out[58] that R parity is not the only discrete symmetry that forbids B- and L-violating dimension-four interactions involving particles in the minimal supersymmetric extension of the Standard Model. There is just one other example[58], a Z_3 symmetry called $GBPR_3L_3$, under which the Standard Model supermultiplets have the following transformation properties:

$$(Q,U,D,L,E,H,\bar{H}) \rightarrow (Q,\frac{1}{\alpha}U,\alpha D,\frac{1}{\alpha}L,\frac{1}{\alpha}E,\frac{1}{\alpha}H,\alpha\bar{H}) : \alpha \in Z_3 \tag{3.77}$$

This has the advantage over R parity that it also forbids some B-violating dimension-five interactions allowed by R parity. The cubic and quartic superpotential terms forbidden by $GBPR_3L_3$ are

$$F_V \ni \lambda LLE^c + \lambda' LQD^c + \lambda'' U^c D^c D^c + QQQL + UUDE + QQQH \tag{3.78}$$

after transformation of the H and L superfields to a basis where there is no quadratic $\bar{H}L$ mixing term. The dimension-five interaction terms allowed by $GBPR_3L_3$ are

$$\mathcal{L}_X \ni (QUEH + LL\bar{H}\bar{H} + LH\bar{H}\bar{H})_F + (\bar{H}\bar{H}E^* + \bar{H}^* HE + QUL^* + UD^* E)_D \tag{3.79}$$

The cosmological bounds (3.24) can be applied directly to the $GBPR_3L_3$-violating interactions (3.78) [3] as noted in the Table, and also to the $GBPR_3L_3$-conserving interactions (3.79), for which the lower bounds are

$$M \gtrsim 10^9 \text{ GeV} \tag{3.80}$$

at a temperature of 1 TeV. (In general, $M > (10^{15}/T)^{1/2n}\, T$ GeV where $n \equiv D - 4$ and T is in GeV.) The phenomenology of $GBPR_3L_3$-violation that follows from the upper bounds in the Table echoes that of R-violation which was discussed in the previous subsections, and will not be discussed here. Note that the limits (3.80) are, in general, bounds on a mass scale divided by coupling constants, so must be interpreted with some care.

[3]These bounds were not evaluated by the authors of Ref. [58], who preferred to exploit the loopholes mentioned in paragraph 4 of Section 1.

3.5 Cosmological constraints on $\Delta B = 2$ interactions

a) $\Delta B = 2$ *interactions*

Certain extensions of the Standard Model contain new interactions which violate baryon number by two units. Possibilities include both $\Delta B = 2$, $\Delta L = 0$ interactions which could induce neutron-antineutron oscillations or nucleon-nucleon annihilation[75], and $\Delta B = \Delta L = 2$ interactions which could induce hydrogen-antihydrogen oscillations[76]. The latter type of interaction conserves $B - L$, like the non-perturbative electroweak effects, so they together may not be able to wash out the BAU. This depends on the explicit field content of the operators, so for convenience we will neglect this class of interactions.

The low-energy effects of $\Delta B = 2$, $\Delta L = 0$ interactions may be represented by operators in an effective lagrangian involving only the light fields of the theory. For a theory with the particle content of the minimal Standard Model, the lowest-dimension $\Delta B = 2$, Lorentz and gauge invariant operators are six-quark interactions of dimension nine

$$\mathcal{L}_{SM}^{\Delta B=2} \sim \frac{1}{M^5} \left(qqq\; qqq\right) \tag{3.81}$$

where for a single generation there are four independent operators of different Lorentz, weak isospin, and colour structure.

In the case of theories with the low-energy field content of the minimal supersymmetric extension of the Standard Model, a $\Delta B = 2$, $\Delta L = 0$ interaction first appears as a dimension-seven F-term

$$\mathcal{L}_{SUSY}^{\Delta B=2} \sim \frac{1}{M_s^3} \left(U^c D^c D^c U^c D^c D^c\right)_F \quad . \tag{3.82}$$

To produce an effective six-quark operator, this F-term must have the external squark lines "dressed" by gluino exchange, resulting in an effective interaction of the form

$$\mathcal{L}^{\Delta B=2}_{SUSY\ eff} \sim \frac{\alpha_s{}^2}{\pi}\frac{1}{\tilde{m}^2}\frac{1}{M_s^3}\left(q^c q^c q^c\ q^c q^c q^c\right) \tag{3.83}$$

where $\tilde{m}$ is the gluino mass scale.

b) *Cosmological and laboratory limits*

At present the best laboratory limit on the $\Delta B = 2$, $\Delta L = 0$ interaction comes from Kamiokande limits on the non-observation of nucleon-nucleon annihilations in the ^{16}O nucleus[41]. They quote a lower limit on the inverse annihilation rate in ^{16}O of 3.4×10^{31} years at 90 % C.L., corresponding to a lower limit of 1.2×10^8 sec. at 90 % C.L. for the free neutron-antineutron oscillation time. If we use the estimates for the hadronic matrix element of the six-quark operator (3.81) from Costa and Zwirner[33] or Mohapatra[42], we find that

$$M > 10^{5 \text{ or } 6}\ \text{GeV} \ . \tag{3.84}$$

The same estimate (and limit) would apply to the "dressed" supersymmetric interaction (3.83) giving

$$M_s \gtrsim (100\ \text{GeV}/\tilde{m})^{2/3} \times 10^7\ \text{GeV} \ . \tag{3.85}$$

If we now compare these to our cosmological rate estimate (3.17) and demand that out-of-equilibrium condition (3.24) be respected at temperatures $T < T_{max}$, we see that the direct laboratory limits are several orders of magnitude weaker than the cosmological limits calculated at $T_{max} \sim 10^{12}$ GeV:

$$M > 10^{14.4}\ \text{GeV}, \quad M_s \geq 10^{14.7}\ \text{GeV} \tag{3.86}$$

but stronger than the bounds computed at $T_{max} \sim 1$ TeV:

$$M > 10^4\ \text{GeV}, \quad M_s \geq 10^5 \left(\frac{100\ \text{GeV}}{\tilde{m}}\right)^{2/3}\ \text{GeV} \tag{3.87}$$

So, if neutron-antineutron oscillations were seen in the laboratory close to present experimental limits, our analysis suggests that the baryon asymmetry of the Universe was either produced late (at the EPT?) or as an e_R asymmetry.

CHAPTER FOUR

EQUILIBRIUM CONSTRAINTS

4.1 Equilibrium in the early Universe

If a particle species in the early Universe is interacting sufficiently fast ($\Gamma \gg H$), the timescale on which the temperature is changing is much longer than that on which the interactions take place, so one can pretend that the particle species i has an equilibrium number density:

$$n_i(m_i, \mu_i, T) = \frac{g_i}{2\pi^2} \int_0^\infty \frac{p^2 dp}{e^{(E-\mu)/T} \pm 1} \tag{4.1}$$

where - (+) is for bosons (fermions), g_i is the number of spin states, m_i is the mass and μ_i is the chemical potential ($\mu < m$ for bosons). In the limit $m_i, \mu_i \ll T$, (4.1) can be approximated as

$$n_i \simeq \frac{g_i}{2\pi^2} T^3 \int_0^\infty \frac{(x^2 + ax)(1 + a/x + (a^2 - b^2)/ax^2)dx}{e^x \pm 1} \tag{4.2}$$

where $x = (E - \mu)/T, a = \mu/T, b = m/T$ and one usually neglects the chemical potential and mass terms because they are very small.

The chemical potential of the i^{th} particle species in a thermal bath is defined as

$$\mu_i = \frac{\partial F}{\partial N_i} \tag{4.3}$$

where F is the Helmholtz free energy and N_i is the number of particles 'i' in the bath. Intuitively, μ_i is the energy (in excess of that of the particle) that one must add to the system when one adds a particle, so the anti-particle number density will be (4.1) with the opposite sign for μ. For $\mu, m \ll T$, the particle asymmetry can therefore be

..pproximated as

$$n_i - \bar{n}_i \simeq \frac{2g_i}{\pi^2} T^2 \mu_i \int_0^\infty \frac{x dx}{e^x \pm 1} \simeq \frac{g\mu T^2}{6} \begin{cases} 2 & \text{bosons} \\ 1 & \text{fermions} \end{cases} \tag{4.4}$$

If one neglects the masses of the particles in the early Universe, the baryon and lepton asymmetries of the Universe in the Standard Model and the minimal supersymmetric Standard Model are proportional to sums of chemical potentials (when masses are neglected):

$$\begin{aligned}
B &= B_L + B_R &&\propto \sum_{gen.=i} \frac{N_c}{3}\{(u_L^i + d_L^i) + (u_R^i + d_R^i)\} \\
L &= L_L + L_R &&\propto \sum_{gen=i} \{(e_L^i + \nu_L^i) + e_R^i\} \\
B &= B_L + B_R + B_S &&\propto \sum_{gen.=i} \frac{N_c}{3}\{(u_L^i + d_L^i) + (u_R^i + d_R^i) + 2(U_L^i + D_L^i + U_R^i + D_R^i)\} \\
L &= L_L + L_R + L_S &&\propto \sum_{gen=i} \{(e_L^i + \nu_L^i) + e_R^i + 2(E_L^i + V_L^i + E_R^i)\}
\end{aligned} \tag{4.5}$$

where B_S and L_S are the baryon and lepton number carried by scalars, N_c is the number of colours and the chemical potential for a particle is represented by the symbol for that particle. It will be useful later to have the B and L carried by the fermions divided into LH and RH components, because the sphalerons take $B_L + L_L \rightarrow 0$.

Equation (4.5) is strictly true only when the particles are massless [36, 37] which we assume to be the case above the electroweak phase transition. The asymmetry in relativistic leptons, for example, of mass m_i and chemical potential μ_i would be $\mu_i f(m_i/T)$, so that even if $\sum \mu_i = 0$, (lepton number violating interaction in equilibrium), one would not get $L = 0$. However, since the sphalerons go out of equilibrium at a temperature that is comparable with that of the EPT, this does not substantially affect the bounds computed in the previous chapter. (Except, perhaps in the case of an Affleck and Dine [79] baryogenesis scenario with a reheat temperature below that of the EPT.)

When the number density of a particle species "α" in a thermal bath is constant, it is said to be in chemical equilibrium [78]. If it is kept in chemical equilibrium via an interaction:

$$\sum_{i=1}^{p} n_i I_i \rightarrow \sum_{j=1}^{q} m_j O_j \tag{4.6}$$

where n_i (m_j) is the number of incident (outgoing) particles I_i (O_j) and α is one of the i or j, then demanding that the total derivative of the free energy with respect to N_α [78] be zero gives

$$\sum_l \frac{\partial F}{\partial N_l}\frac{\partial N_l}{\partial N_\alpha} = \sum_{i=1}^{p} n_i\mu_i - \sum_{j=1}^{q} m_j\mu_j = 0 \ . \tag{4.7}$$

So when a reaction is in equilibrium, the sum of the chemical potentials of the ingoing particles is equal to that of the out-going and every such reaction constrains the chemical potentials to satisfy a homogeneous equation. Solving for the potentials is a simple exercise in matrix algebra; if there are m interactions in equilibrium, involving n particles, then one can write this as the matrix equation

$$M\vec{\mu} = 0 \tag{4.8}$$

where M is an $m \times n$ matrix. If $m \geq n$, then $\vec{\mu} = (\mu_1, \mu_2, ..., \mu_n)^T = 0$, and the interactions take all the particle asymmetries to zero. If instead $n - m = p > 0$, then there are p symmetries or possible combinations of chemical potentials that sum to a constant = a conserved quantum number. (These are not neccessarily symmetries of the lagrangian; there could be out-of-equilibrium interactions that violate them.)

4.2 Chemical potentials in the Standard Model

Above the electroweak phase transition, the gauge bosons are massless and should have zero chemical potential (unless there is a condensate, but this would break the gauge symmetry). One can show [6] that requiring the thermal bath to have $Q_3 = 0$ (where Q_3 is the charge associated with the diagonal SU(2) generator) implies that the chemical potential of the W's is zero. A similar argument should apply to the gluons, so one can conclude that all particles in a gauge multiplet have the same chemical potential. If one just considers gauge interactions, there are 16 independent chemical potentials: $h, \ell^i_L, q^i_L, e^i_R, u^i_R, d^i_R$, where i is a flavour index running from 1 to 3. Requiring that the total electric charge, or hypercharge, be zero gives another constraint:

$$Q_{em} = \frac{N_g N_c}{3}(q_L + 2u_R - d_R) - \sum_i \ell^i_L - \sum_i e^i_R + h = 0 \tag{4.9}$$

which can be used to solve for h as a function of chemical potentials carrying B or L.

Including the Higgs interactions implies

$$\begin{aligned} -h + \ell^i_L - e^i_R &= 0 \\ -h + q^i_L - d^j_R &= 0 \\ h + q^i_L - u^j_R &= 0 \end{aligned} \tag{4.10}$$

The lepton-Higgs interactions leave one independent lepton chemical potential per generation. Since the quark-Higgs couplings are not generation diagonal, there will be one free quark potential. The free chemical potentials correspond to the four global symmetries of the SM Lagrangian: L_i, B. Note, however, that unless $h = 0$, the asymmetries stored in the left and right-handed particles will not be equal.

The electroweak $B+L$ violating processes require

$$3N_g q_L + \sum_i \ell_L^i = 0 \qquad (4.11)$$

leaving the three global symmetries $B/3 - L^i$. Now suppose that all the above Standard Model processes, and the sphalerons, are in thermal equilibrium; if any interaction violating all the L^i or $B/3 - L^i$ comes into thermal equilibrium, it will provide three more homogeneous equations for (4.8), and force every chemical potential to zero.So any L or $B-L$ violating interaction with generic generation non-diagonal couplings will take the baryon asymmetry to zero.

However, one does not have to force all the chemical potentials to zero to wipe out the BAU; one just needs

$$B = N_g(2q_L + u_R + d_R) = 4N_g q_L = 0 \qquad (4.12)$$

so a "beyond the Standard Model" B violating interaction that takes q_L to zero will remove any existing baryon asymmetry. This is independent of non-perturbative electroweak effects (equation 4.11). However, if the sphalerons are in thermal equilibrium, any B violating interaction will wipe out the BAU, not just those that take q_L to zero. (Algebraically, this is because one can express h as a function of q_L alone using (4.11), so that q_L or u_R or $d_R = 0$ implies $h = 0$ and $u_R = d_R = q_L = 0$.) It is also worth noting that L violating interactions do not need to take each family asymmetry to zero; one only needs $L^1 + L^2 + L^3 \rightarrow 0$.

It has just been shown that any B, L, or $B-L$ violating interaction in thermal equilibrium with the sphalerons would wipe out a pre-existing BAU, which is the premise we used to calculate bounds on higher dimensional SM operators in the previous chapter. Such constraints get better for high temperatures, so one would like to know the maximum temperature at which they are applicable ... which unfortunately is rather low.

The gauge interactions are quite strong, so should stay in thermal equilibrium up to higher temperatures than the "sphalerons". The Higgs Yukawa couplings are in some cases very small, and taking the rate for $h \rightarrow e_L e_R$ to be

$$\Gamma \sim 10^{-2}\frac{h_{ee}^2 m_h^2}{T} \tag{4.13}$$

then $\Gamma < H$ above $T \sim 1$ TeV. One can easily calculate the temperature at which the other Higgs couplings go out of equilibrium, but this is not the point. If the RH electrons are not in chemical equilibrium with the rest of the SM particles, one could in principle store an asymmetry in them, allow arbitrary B and/or L interactions to be in thermal equilibrium with the sphalerons down to $T > 1$ TeV, and then transfer the e_R asymmetry to the baryons via the higgs coupling and the sphalerons. The maximum temperature at which our constraints apply therefore depends on the baryogenesis model. If the asymmetry is created in the LH fermions, we can use $T_{max} \approx 10^{12}$ GeV, and get very strong limits. If it is in some RH fermion, then the bounds are weaker.

4.3 The Minimal Supersymmetric Standard Model

In the MSSM, the scalar partners of the leptons and quarks carry B and L (see equations 4.5). In an exactly supersymmetric theory, they would be massless above the electroweak phase transition, and could only carry an asymmetry in a condensate. However, we expect supersymmetry to be broken by explicit mass terms for the superpartners because they have not been observed, so the squarks and sleptons can have chemical potentials.

As in the SM, the gauge interactions will force all members of a gauge multiplet to have the same chemical potential (providing the gauge symmetry is unbroken).

The quartic scalar field couplings from the F and D terms of the supersymmetric lagrangian do not bring particles into chemical equilibrium (no constraints on the chemical potentials) because they always involve particle anti-particle pairs. The D term Yukawas, involving a gaugino and the fermion and scalar components of a superfield imply:

$$\phi + \lambda + \psi = 0 \tag{4.14}$$

where ϕ, ψ and λ are the chemical potentials for the matter scalar and fermion and the gaugino respectively. As a particlular example of this, one has

$$\begin{aligned} Q_L + \tilde{g} + q_L &= 0 \\ Q_L + \tilde{W} + q_L &= 0 \\ Q_L + \tilde{B} + q_L &= 0 \end{aligned} \tag{4.15}$$

where Q_L is the "left-handed" scalar quark doublet, $\tilde{g}$ the gluino, $\tilde{W}$ the wino and $\tilde{B}$ is the fermion partner of the hypercharge gauge boson. All the gauginos therefore have the same chemical potential, and one can calculate the chemical potential of any scalar knowing that of the gauginos and the fermion superpartner.

Since we are working above the EPT, we assume that the gauge bosons have zero chemical potential. After including, as above, constraints on other particles due to the gauge interactions being in chemical equilibrium one is left with eighteen free chemical potentials which can be taken to belong to the fermions: λ (the gaugino chemical potential), ℓ^i_L, e^{ci}_R, q^i_L, u^{ci}_R, d^{ci}_R, h and $\bar{h}$, where i is a generation index, fermions are lower case letters and scalars are capitals (so h and $\bar{h}$ are the superpartners of the Higgs). Demanding that the electric charge of the thermal bath be zero gives one more constraint.

The Higgs mixing term $\varepsilon H\bar{H}$ will force H to have the opposite chemical potential from $\bar{H}$. Each trilinear term in the superpotential contributes three Yukawa interactions to the lagrangian, and these will share the same equation of chemical equilibrium, given the relation between scalars and fermions in (4.15). Using $h = -\bar{h}$, these will be

$$
\begin{aligned}
\ell^i_L - h - (e^{ci}_R + \lambda) &= 0 \\
q^i_L - h - (d^{cj}_R + \lambda) &= 0 \\
q^i_L + h - (u^{cj}_R + \lambda) &= 0
\end{aligned}
\tag{4.16}
$$

There are now five chemical potentials left, which can be taken to be ℓ^i_L, q_L and λ.

If the sphalerons are in thermal equilibrium, there is one more constraint. However, it is not the same as in the Standard Model, because gauge field configurations of finite winding number couple to the axial current, and there are chiral fermions in the MSSM that do not carry baryon or lepton number.The appropriate equation is [13]

$$
[\text{illegible}]_g q_L + \sum_i \ell^i_L + (h + \bar{h}) + 4\lambda = 0 \quad . \tag{4.17}
$$

This allows us to compute λ in terms of B_L and L_L, but imposes no constraints on B or L unless we already know λ. One could therefore have B or L violating interactions in thermal equilibrium with the sphalerons without wiping out the BAU. However, one expects SUSY breaking majorana masses for the gauginos (we have not seen gluinos), which would force $\lambda \rightarrow 0$, so that (4.17) would reduce to the same equation as for the Standard Model, and all the arguments of the previous section would apply.

It has been shown [13] that at temperatures sufficiently high that the $\varepsilon H\bar{H}$ Higgs mixing term and the SUSY breaking mass terms are out of thermal equilibrium, the MSSM has extra anomalous global symmetries. One then finds [13] that the sphalerons take $B + L +$ 'the conserved charge of one of these symmetries' to zero, so that in the presence of B, L or $B - L$ violation, the BAU is not washed out. However, the right-handed electron drops out of chemical equilibrium before the supersymmetry-breaking mass terms, so this argument will not be reviewed in detail here.

4.4 Protecting the BAU

In this thesis, I have used B to represent both the quantum number and the asymmetry, which naturally leads to some confusion. In the first sense, B is a quantum number: +1/3 for quarks of all colours, chiralities and flavours. In the second sense, B is the sum of the particle-anti-particle asymmetries for all the baryon-number carrying particles (see equation 4.5). The confusion arises because interactions that violate B in the first sense (B for the rest of this section) do not necessarily take B in the second sense (b for the rest of the section) to zero. For instance, as pointed out by Dreiner and Ross, [36], the $U^cD^cD^c$ superpotential term in the MSSM would not wash out the BAU in the absence of sphalerons because it gives

$$u_R + 2d_R = 3q_L - h = 0 \tag{4.18}$$

which does not imply $b(= 12q_L) = 0$ unless h is known as a function of q_L (as is the case when sphalerons are in equilibrium). This interaction allows two d_R quarks to turn into an anti-U_R squark, so in equilibrium the excess of $\bar{U}_R$ will be twice that of the d_Rs. This does not force the BAU to zero because the chiral structure of the SM puts different fractions of the baryon asymmetry into the u_R, d_R and q_L.

For the same reason, although the sphalerons violate $B + L$, they do not take $b + l$ to zero; one can easily see, from (4.9), (4.5) and (4.10) that (4.11) implies

$$b + \frac{52}{93}l = 0 \tag{4.19}$$

If all the Standard Model interactions and the "sphalerons" are in chemical equilibrium, any interaction forcing b or l to zero will wipe out the BAU. Hence the constraints of chapters 2 and 3. However, although any B violating interaction takes b to zero in this case, it is not true that any L violating interaction will take l to zero. This is because the SM conserves each lepton family number individually. So if one of the lepton family numbers is accidentally conserved in a $B - L$ violating interaction, l will not go to zero, and the baryon asymmetry will survive [5]. There is no particular reason to assume that an L-violating interaction would respect one of the lepton family numbers; however, if it was incapable of taking l to zero, none of the constraints calculated in this thesis would apply.

It has been previously mentioned that the smallness of the Higgs coupling between left and right handed fermions makes our constraints on higher dimensional operators weaker than one might initially think. This is a variation on the observation that an accidentally conserved quantum number can protect the BAU: if a right-handed particle is out of chemical equilibrium with respect to the rest of the SM particles, it behaves as if it carried an accidentally conserved quantum number, and any asymmetry stored in it is effectively conserved until the particle comes into chemical equilibrium. It is possible that our constraints still apply at $T > 1$ TeV (when the e_R go out of chemical equilibrium), but this becomes a model dependent statement. For instance, if the asymmetry is produced in the LH fermions, the constraints apply up to $\sim 10^{12}$ GeV (10^9 GeV in the MSSM [13]) and are very strong; on the other hand, if the asymmetry is in the right-handed electrons, $B - L$-violating interactions can be in equilibrium down to temperatures of order 1 TeV.

CHAPTER FIVE

CONCLUSION

In chapter two we presented order-of-magnitude constraints on B- and L-violating extensions of the Standard Model that follow from requiring that a primordial cosmological baryon asymmetry survive despite the strong non-perturbative electroweak interactions that violate $B+L$ and conserve $B-L$. In chapter three, we examined this scenario more carefully, giving quantitative rate estimates for different interactions and obtained generic bounds on the corresponding couplings. We then made applications of these bounds to explicit and spontaneous lepton number violation, and to the violation of R [68] (and $GBPR_3L_3$) [58] parity in supersymmetric models. We also discussed cosmological constraints on $\Delta B = 2$ interactions that can be independent of the existence of non-perturbative electroweak interactions.

As seen in the tables, our bounds are in general much more stringent than those coming from present laboratory experiments. Thus, for example, if our arguments apply, $\not{R}$-violating decays of the Z^0 or of the LSP should be unobservable, and the only possible observable signature of $\not{R}$-violation would be a small window for the LSP to be a massive charged or strongly-interacting particle which would live long enough to appear stable in any laboratory experiment, but with a lifetime $\lesssim 10^4$ s. Cosmological arguments also suggest that $n - \bar{n}$ oscillations or $N - N$ annihilations in nuclei should be unobservable.

As we pointed out in the introduction to chapter 3, there are several loopholes in our cosmological arguments. For example, generating the cosmological baryon asymmetry at some temperature $T < T_c \sim 100$ GeV seems to require new physics beyond the Standard Model at the electroweak scale (an extra Higgs doublet ? additional CP-violation ?). Spontaneous R-violation is also likely to require an extension of the electroweak sector of the Standard Model. In fact, the only way of avoiding our constraints without introducing new physics appears to be to protect the BAU with a conserved (or effectively conserved) quantum number.

One could therefore conclude that the observation of a large Majorana neutrino mass, *or* of R- or $GBPR_3L_3$-violation, *or* any other interaction violating our constraints could be taken as evidence that *other* new physics could show up in accelerator experiments.

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APPENDIX A

TABLES

Interaction	D	T = 100 GeV	lab limit
$\mu^2 \tilde{l} h$	2	$\mu < .1$ GeV	$\mu < 1 - 100$ GeV
$m_3^s \tilde{q} \tilde{d}^c \tilde{l}$	3	$m_3^s < 10^{-5}$ GeV	none ?
$m_3^s \tilde{l} \tilde{e}^c \tilde{l}$	3	$m_3^s < 10^{-5}$ GeV	none ?
$m_3^s \tilde{u}^c \tilde{d}^c \tilde{d}^c$	3	$m_3^s < 10^{-5}$ GeV	none ?
$m_3^f l \tilde{h}$	3	$m_3^f < 3 \times 10^{-5}$ GeV	$m_3^f < 0.1 \frac{\varepsilon}{30 \text{ GeV}}$

Table A.1: bounds on B or L violating soft ($D \leq 3$) supersymmetry-breaking interactions

Interaction	D	T = 100 GeV	lab limit
$m(LH)_F$	3	$m < 3 \times 10^{-5}$ GeV	$m < 0.1 \frac{\varepsilon}{30 \text{ GeV}}$
$\lambda_1(LLE^c)_F$	4	$\lambda < 10^{-7}$	$\lambda_1^{113} < 2 \times 10^{-3} \frac{\tilde{m}}{100 \text{ GeV}}^{1/2}$
$\lambda_2(QD^cL)_F$	4	$\lambda < 10^{-7}$	$\lambda_2^{113} < 10^{-3} \frac{\tilde{m}}{100 \text{ GeV}}^{1/2}$
$\lambda_3(U^cD^cD^c)_F$	4	$\lambda < 10^{-7}$	$\lambda_3^{11k} < 2 \times 10^{-7} \frac{\tilde{m}}{100 \text{ GeV}}^{5/2}$

Table A.2: bounds on renormalizable ($D \leq 4$) supersymmetric B or L violating interactions

Interaction	D	T = 100 GeV	10^5 GeV	10^7 GeV	10^9 GeV	10^{12} GeV	lab limit
$llhh$	5	10^8 GeV	7×10^9 GeV	7×10^{10} GeV	7×10^{11} GeV	2×10^{13} GeV	$10^{11} - 10^{15}$ GeV
$qqql^ch$	7	10^4 GeV	4×10^6 GeV	2×10^8 GeV	10^{10} GeV	3×10^{12} GeV	$10^{10} - 10^{11}$ GeV
$qqql^c\mathcal{D}$	7	10^4 GeV	4×10^6 GeV	2×10^8 GeV	10^{10} GeV	3×10^{12} GeV	$10^8 - 10^9$ GeV
$qqqqqq$	9	2×10^3 GeV	10^6 GeV	6×10^7 GeV	4×10^9 GeV	2×10^{12} GeV	$10^5 - 10^6$ GeV
$qqqlll$	9	2×10^3 GeV	10^6 GeV	6×10^7 GeV	4×10^9 GeV	2×10^{12} GeV	none ?
$qqql^cl^cl^ch$	10	10^3 GeV	6×10^5 GeV	4×10^7 GeV	3×10^9 GeV	2×10^{12} GeV	$10^4 - 10^5$ GeV
$qqqlllhh$	11	8×10^2 GeV	5×10^5 GeV	4×10^7 GeV	3×10^9 GeV	2×10^{12} GeV	10^4 GeV

Table 1: bounds on $D > 4$ B or L violating Standard Model interactions

Interaction	D	T = 100 GeV	10^5 GeV	10^7 GeV	10^9 GeV	10^{12} GeV	lab limit
$(LLHH)_F$	5	10^8 GeV	7×10^9 GeV	7×10^{10} GeV	7×10^{11} GeV	2×10^{13} GeV	$10^{11}-10^{15}$ GeV
$(L\bar{H}HH)_F$	5	10^8 GeV	7×10^9 GeV	7×10^{10} GeV	7×10^{11} GeV	2×10^{13} GeV	none?
$(QQQ\bar{H})_F$	5	10^8 GeV	7×10^9 GeV	7×10^{10} GeV	7×10^{11} GeV	2×10^{13} GeV	none?
$(D^{c*}D^{c*}D^{c*}E^c)_D$	6	10^5 GeV	3×10^7 GeV	8×10^8 GeV	3×10^{10} GeV	5×10^{12} GeV	10^{11} GeV
$(D^{c*}D^{c*}QL^*)_D$	6	10^5 GeV	3×10^7 GeV	8×10^8 GeV	3×10^{10} GeV	5×10^{12} GeV	10^{11} GeV
$(D^{c*}D^{c*}D^{c*}L^*\bar{H}^*)_F$	6	10^5 GeV	3×10^7 GeV	8×10^8 GeV	3×10^{10} GeV	5×10^{12} GeV	10^{11} GeV
$(D^{c*}D^{c*}U^{c*}L^*H^*)_F$	6	10^5 GeV	3×10^7 GeV	8×10^8 GeV	3×10^{10} GeV	5×10^{12} GeV	10^{11} GeV
$(U^cD^cD^cU^cD^cD^c)_F$	7	10^4 GeV	4×10^6 GeV	2×10^8 GeV	10^{10} GeV	3×10^{12} GeV	$10^8\frac{100\text{ GeV}}{\tilde{m}}$ GeV

Table 2: bounds on $D > 4$ B or L violating supersymmetric operators

APPENDIX B

FIGURES

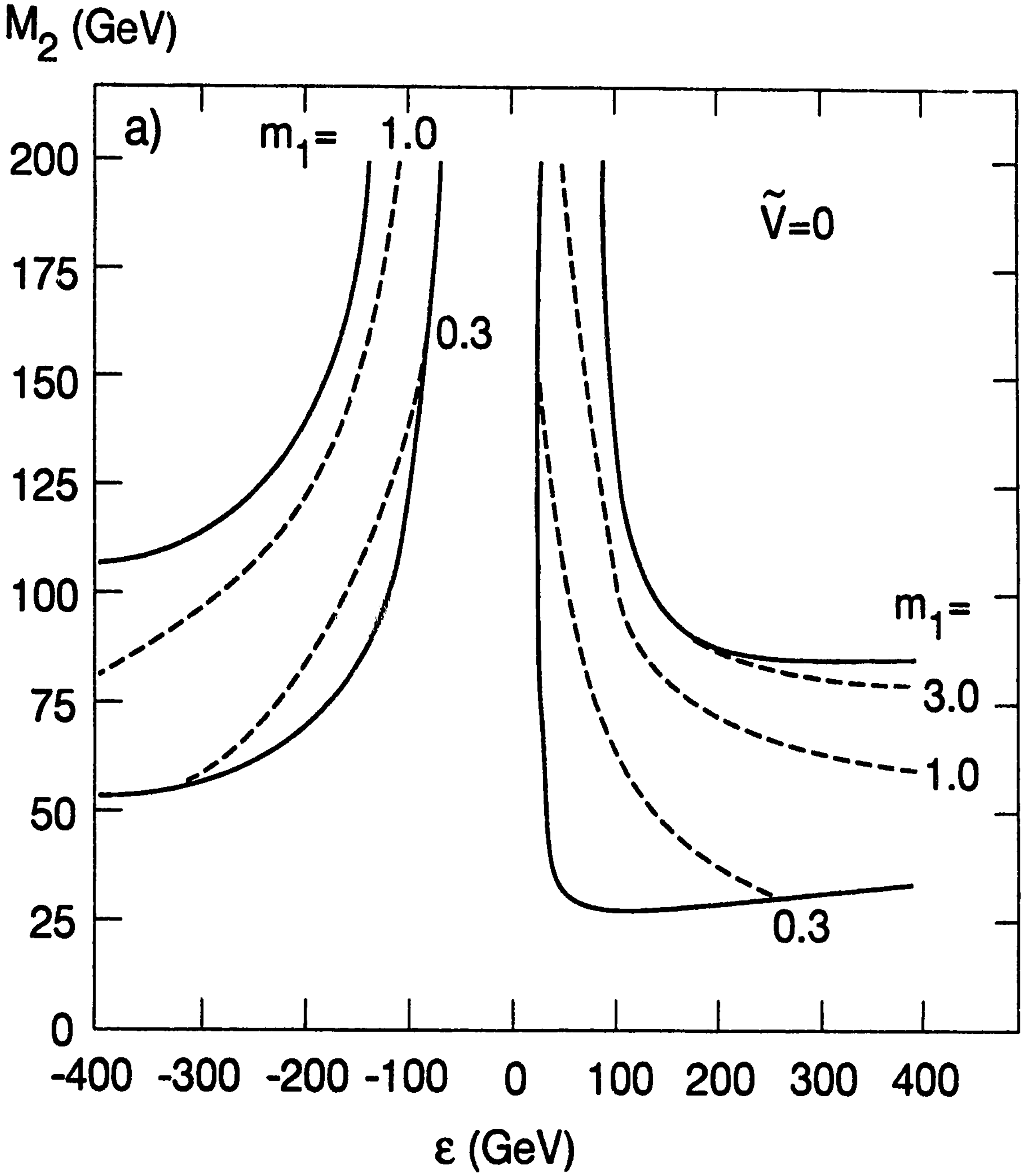

Figure B.1: The R_P violating branching ratio $\mathrm{BR}(Z_O \to \tau^+\chi_-)$ in the (ϵM_2) plane with $\tilde{V} = 0$. The dashed lines are contours of M_1 (in GeV) that would give 10 events at LEP with 10^7 Z_0s. The solid lines correspond to $M_\chi = 40$ (90) GeV.

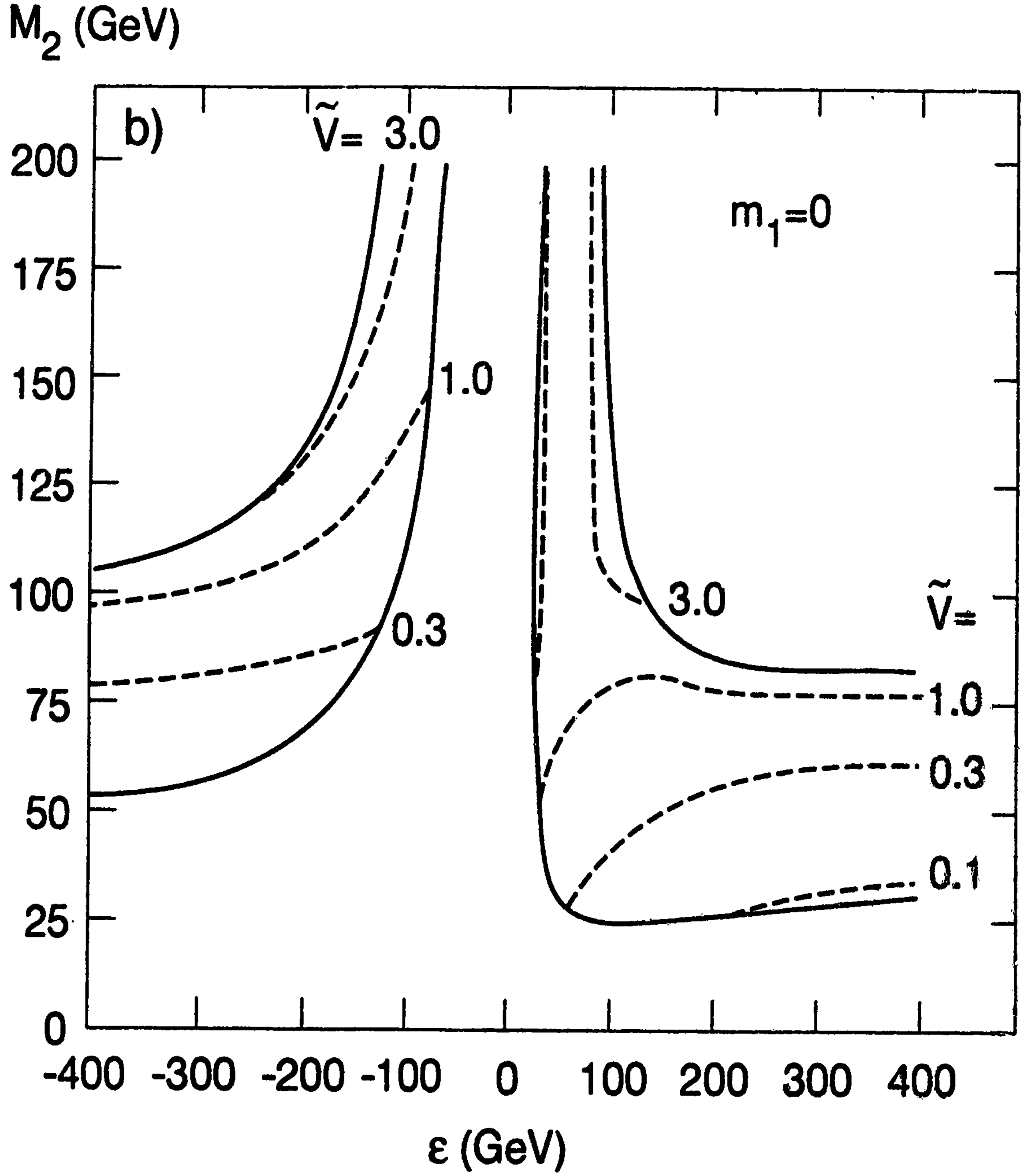

Figure B.2: The R_P violating branching ratio BR($Z_O \to \tau^+\chi_-$) in the (ϵM_2) plane with $M_1 = 0$. The dashed lines are contours of $\tilde{V}$ (in GeV) that would give 10 events at LEP with 10^7 Z_0s. The solid lines correspond to $M_\chi = 40$ (90) GeV.

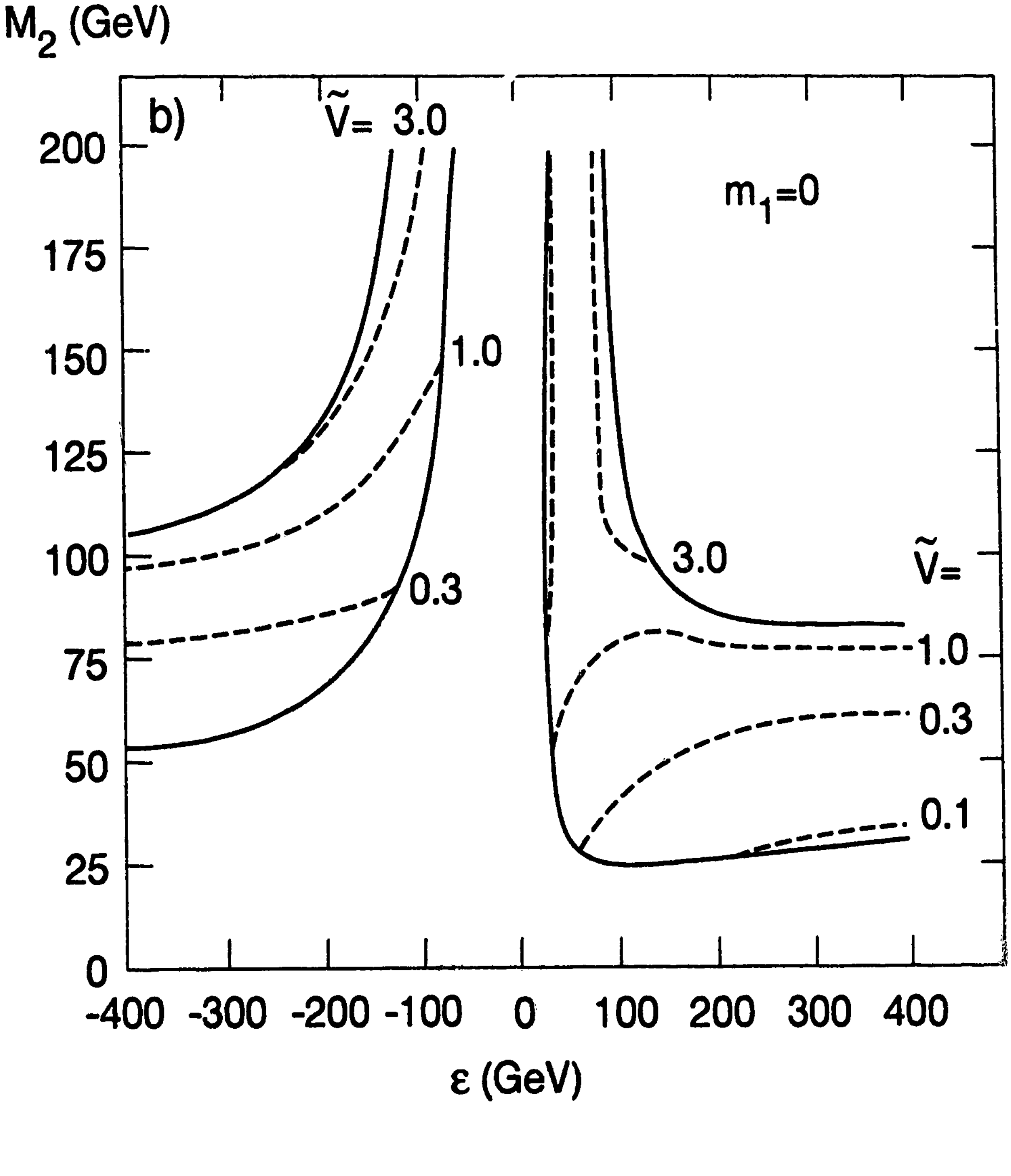

Figure B.3: The R_P violating branching ratio BR($Z_O \to \tau^+\chi_-$) in the (ϵM_2) plane. The dashed lines are contours of $M_1 = \tilde{V}$ (in GeV) that would give 10 events at LEP with 10^7 Z_0s. The solid lines correspond to $M_\chi = 40$ (90) GeV