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Beyond the Supersymmetric Standard Model

L.J. Hall

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BEYOND THE SUPERSYMMETRIC STANDARD MODEL *

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I. INTRODUCTION

These lectures are neither a broad overview of the huge field given by the title, nor an exhaustive treatment of a more narrowly defined topic. Instead I have chose to discuss a few ideas which I find exciting. In an attempt to make the lectures stimulating, these ideas are radically different from conventional approaches to supersymmetric unification and they are topics of current research. These ideas include the possibility of baryon number violation at the weak scale and an alternative primordial nucleosynthesis scheme arising from the decay of gravitinos. Both of these ideas involve the breaking of R parity and consequently lead to unusual and spectacular signatures for supersymmetry at e^+e^- and hadron colliders.

I have not emphasized the formalism or the calculational tools. The emphasis is entirely on the physics ideas. I hope that the reader finds the physics sufficiently exciting to learn the necessary tools from elsewhere. These lectures inadequately cover the basics of grand unification, the minimal supersymmetric standard model and conventional supersymmetric unification. Fortunately, these subjects are already well treated in previous TASI lectures in references 1, 2 and 3 respectively.

II. SELECTED FEATURES OF SUPERSYMMETRIC MODELS

In this section I do three things: I briefly define the minimal low energy supergravity (MLES) model, I discuss a few of its features which we will later need, and I mention renormalization group scaling and flavor physics.

II.1 The Minimal Low Energy Supergravity Model

We must at least define the MLES model[2] which we will be extending. This is a supersymmetric $SU(3) \times SU(2) \times U(1)$ gauge theory with three 15-plets of chiral superfields for "matter"

$$Q(3,2,1/6) \quad U^c(\bar{3},1,-2/3) \quad D^c(\bar{3},1,1/3) \quad L(1,2,-1/2) \quad E^c(1,1,1)$$

and two for "Higgs"

$$H(1,2,1/2) \quad H'(1,2,-1/2).$$

Supersymmetry itself does not allow for a distinction between matter and Higgs fields, so we impose one by hand: we require the MLES model to be invariant under matter parity under which the matter superfields change sign but the Higgs superfields do not. The most general gauge invariant, renormalizable superpotential is then

$$f = Q\lambda_U U^c H + Q\lambda_D D^c H' + L\lambda_E E^c H' + \mu H H' \qquad (1)$$

where λ_U, λ_D and λ_E are 3×3 matrices in generation space and μ is a dimensionful parameter which ensures that the theory does not possess a Peccei-Quinn symmetry.

A useful way of remembering how to get the vertices of the supersymmetric interactions of Equation (1) in terms of component fields is to write down the usual Yukawa couplings and replace external lines in pairs by their superpartners as in Figure 1.1(a). Finally scalar trilinear and quartic interactions are generated by differentiating f, considered as a function of the scalar components of the superfields A_i, as shown in Figure 1.1(b):

The model that I am describing, despite its name, doesn't have anything to do with supergravity. However, a crucial aspect of the model is the structure of the soft supersymmetry breaking operators. These operators can be obtained in a very plausible fashion from supergravity theories (for a detailed discussion and review of the possibilities see reference 4), but we will not need to know anything about supergravity. In the simplest scheme there are four types of soft operators:

$$m^2 A_i^* A_i = m^2(H^* H + \tilde{q}^* \tilde{q} + \cdots) \tag{2a}$$

$$Bm[f_2]_A = Bm\mu H H' \tag{2b}$$

$$Am[f_3]_A = Am\tilde{q}\lambda_U \bar{u}^c H + \cdots \tag{2c}$$

$$-\frac{\tilde{m}}{2}(\tilde{g}\tilde{g} + \tilde{\omega}\tilde{\omega} + \tilde{b}\tilde{b}) \tag{2d}$$

$[f_{2,3}]_A$ are the bilinear, trilinear parts of the superpotential as functions of the scalar components of the superfields A_i. A and B are complex constants with magnitudes of order unity, and $\tilde{g}, \tilde{\omega}, \tilde{b}$ are the gaugino fields for $SU(3), SU(2), U(1)$ gauge groups. Note that H sometimes refers to a superfield and sometimes to its scalar component. Superpartners are differentiated from the standard model particles by a tilde.

Is the MLES just an irrelevant extension of the standard model which has introduced five new parameters and a host of new particles for naught? Experiment must decide. It is theoretically attractive because, unlike the standard model, the theory has no quadratic divergences. All parameters scale according to well behaved renormalization group Equations (RGE). If we write a supersymmetric theory with a supersymmetry breaking scale much above the weak scale, then on integrating out the superpartners we will recover the standard model as a low energy effective theory with its quadratically divergent Higgs mass. To prevent this, the supersymmetry breaking and weak scales should be comparable. The MLES model incorporates this

automatically: if m and $\tilde{m}$ are made very large the Higgs boson will decouple, hence m and $\tilde{m}$ cannot be made larger than the weak scale v: m and $\tilde{m}$ are taken to be $0(v)$. Similarly if $\mu \gg v$ the Higgs decouples. The most puzzling feature of MLES is why the mass parameters in the supersymmetric and supersymmetry breaking vertices have comparable sizes.

II.2 The Electric Dipole Moment of the Neutron

The MLES model differs from the standard model in many of its predictions. This is well illustrated by one particularly striking case which will be of importance to us later: the neutron electric dipole moment, d_n. This parameter violates CP and is known to be less than 10^{-24}e cm.[5] Ignoring the strong CP parameter, the standard model violates CP only through the single parameter δ which appears in the Kobayashi-Maskawa matrix. This leads to $d_n \simeq 10^{-31}$ e cm. via a multiloop diagram.

The pattern of CP violation is much richer in the MLES model: there are two additional phases which can be taken to be those of the A and B parameters ϕ_A and ϕ_B.[6] Each of these gives rise to a contribution to d_n at one loop;[7),8),6)] that coming from ϕ_A is shown in Figure 2. At low energies this diagram yields a local operator in the effective Lagrangian, with a coefficient which we crudely estimate

$$\mathcal{L}_{eff} = \left(\frac{e g_s^2}{16\pi^2}\,\phi_A A\,\frac{m m_u \tilde{m}}{m^4 + m^2\tilde{m}^2}\right) \bar{u}\sigma^{\mu\nu}\gamma_5 u F_{\mu\nu}. \tag{3}$$

Here e and g_s are the electromagnetic and strong gauge couplings, and the denominator has been chosen to give the correct behavior whether $m > \tilde{m}$ or vice versa. The coefficient is approximately d_n, which for $m = \tilde{m}$ is

$$d_n \simeq 10^{-23}\ \text{e cm}\ A\phi_A \left(\frac{300\ \text{GeV}}{m}\right)^2 \tag{4}$$

a result up to nine orders of magnitude larger than in the standard model. This is very encouraging for future searches for a non-zero d_n, unless we find a reason why ϕ_A should be very small. The present experimental bound already imposes a non-trivial constraint on the MLES model. For $m, \tilde{m} \leq 100$ GeV, we need to understand why $A\phi_A$ should be much less than unity. If we cannot understand this, perhaps it is an indication that squarks and gluinos may be nearer a TeV in mass.

II.3 R Parity

In the MLES theory all component vertices have an even number of superpartner fields. This means the theory possesses a symmetry under which the sign of these

fields is reversed: R parity. It has the consequence that the lightest superpartner (LSP) is stable.

R parity is a discrete subgroup of a continuous $U(1)$ R symmetry which rotates the coordinate θ of superspace. The superfields can be expanded in terms of θ and component fields, for example

$$Q = \tilde{q} + \theta q + \cdots, \tag{5a}$$

$$H = H + \theta\tilde{H} + \cdots, \tag{5b}$$

$$Z = \cdots + \theta\sigma^{\mu}\bar{\theta}Z_{\mu} + \bar{\theta}\bar{\theta}\theta\tilde{Z} + \cdots \tag{5c}$$

You can now check very simply that R parity in MLES, which reverses the sign of θ and all superpartner component fields, is exactly the same thing as matter parity, which reverses the sign of matter superfields, but not Higgs or vector superfields.

R parity plays a central role in the phenomenology of the MLES model, and largely determines the nature of its experimental signatures. This is for two reasons: i) R parity implies that direct production of superpartners will occur in pairs (although one could be off-shell). ii) Once a superpartner has been made you can never get rid of it (except for the possibility that it might come across another superpartner to annihilate). This is important cosmologically since relic superpartners from the big bang will decay to products which include the LSP, and since the LSP is stable it could be the dark matter. To avoid cosmological problems the LSP should therefore be neutral: a neutralino or sneutrino. The stability of the LSP is also crucial for lab. searches for supersymmetry. Once produced in a high energy collision, a neutral, stable LSP will escape the apparatus and leave a missing energy signature. The vast majority of searches for supersymmetry, and limits on superpartner masses, have used this signature [9].

Since the conservation of R parity plays such a central role in present thinking in supersymmetric models, a central theme of this set of lectures will be to challenge this standard viewpoint and to explore possibilities for R parity breaking. Such breaking can be explicit [10] or spontaneous [11] at the weak scale, or may even be restricted to higher dimension operators [12]. In this set of lectures I will restrict myself to discussing minimal field content (*i.e.* same as in MLES) models with R violation explicit at the weak scale. The four logical possibilities are shown in Table I. The only possible

TABLE I

	Excluded	MLES	$\Delta L \neq 0$	$\Delta B \neq 0$
LLE^c , QD^cL	√	X	√	X
$U^cD^cD^c$	√	X	X	√

gauge invariant, supersymmetry, R violating operators in these models are those which violate lepton number

$$f_{\Delta L} = \lambda_1 LLE^c + \lambda_2 QD^cL + \mu' LH \tag{6}$$

and those which violate baryon number

$$f_{\Delta B} = \lambda_3 U^cD^cD^c. \tag{7}$$

Most theorists would probably have opted for the model in the first column of Table I, since it contains all possible gauge invariant interactions, and this seems most natural. However, this possibility is excluded because the proton decays with a weak decay rate. The next simplest version is to assume that neither is present; this produces the MLES model shown in the second column. One reason for doing this is that grand unification treats quarks and leptons on equal footings and hence suggests that either both or neither type of new coupling should appear. This argument is not quite correct (there is a counterexample in reference 10) and in any case we need not believe in conventional grand unification. Hence in these lectures we will often consider the "$\Delta L \neq 0$" model or the "$\Delta B \neq 0$" model of columns 3 and 4.

It is worth pointing out that it is remarkable how large the couplings $\lambda_{1,2,3}$ can be without conflicting with data. In either model it is possible to have some of these couplings 0(1), while others must be quite small, $0(10^{-6})$ [10,13]. In particular it is important to point out that the proton is stable in the "$\Delta B \neq 0$" model. This is because, from angular momentum arguments, the proton must decay to a final state containing an odd number of leptons (we take all superpartners heavier than the proton), and this is forbidden by lepton number conservation. The most dangerous processes are then heavy "*stable*" nucleus decay, and neutron oscillations [13].

II.4 Renormalization Group Scaling of Parameters

The renormalization group (RG) scaling of parameters in supersymmetric models is much more elegant and contains much more physics than in the non-supersymmetric case. While the scaling of gauge couplings α_i and Yukawa couplings $\lambda_{U,D,E}$ is similar, there is also the scaling of the parameter μ, and, most significantly of all, the

soft parameters: A, B, m_i (the scalar mass parameters) and $m_{\tilde{a}}$ (the gaugino mass parameters).

The initial conditions to the RG equations come from two assumptions

i) Grand unification. If we want to impose a conventional grand unification then the superfields have a higher group multiplet structure. For example, in $SU(5)$:

$$\overline{F} = \begin{pmatrix} D^c \\ L \end{pmatrix} \quad T = \begin{pmatrix} U^c & Q \\ -Q & E^c \end{pmatrix} \quad H_{1,2} = \begin{pmatrix} H_3 \\ H \end{pmatrix}, \begin{pmatrix} H_3' \\ H' \end{pmatrix}$$

are five dimensional tensors, with color acting on the first three indices and weak $SU(2)$ on the last two. Above the unification scale, the boundary condition must be:

$$\alpha_3 = \alpha_2 = \alpha_1 = \alpha_G, \tag{8a}$$

$$m_{d^c} = m_l = m_{\bar{F}}, \tag{8b}$$

$$m_{\tilde{u}^c} = m_{\tilde{q}} = m_{\tilde{e}^c} = m_T \tag{8c}$$

$$m_{\tilde{g}} = m_{\tilde{\gamma}} = m_{\tilde{q}} = \tilde{m}. \tag{8d}$$

A larger symmetry than $SU(5)$ would be required to relate $m_{\bar{F}}$ to m_T.

ii) Supergravity. The idea that the soft operators are induced from a sector of a supergravity theory leads to additional boundary conditions at the Planck scale [14,4]. These are

$$m_{\bar{F}} = m_T = m, \tag{9a}$$

$$(\text{Trilinears}) = A\, m[f_3]_A, \tag{9b}$$

$$(\text{Bilinears}) = B\, m[f_2]_A. \tag{9c}$$

The latter two relations are modified at low energies by RG scaling; the flavor structure of the trilinear scalar couplings will not be the same as that of the Yukawa couplings.

After these preliminaries, we discuss a few salient features of the RG scaling. More complete presentations are given in the references.

(a) Gauge couplings g_a

A comparison between the grand unification of gauge couplings in ordinary [15] and SUSY [16] theories is given schematically in Figure 3. Above the grand unification scale there is a single gauge coupling. At energies low enough such that the non-degeneracies of irreducible representations of $SU(5)$ are apparent the gauge couplings

diverge. This is because the effects of heavy particles decouple; they do not contribute to the β function coefficient b_a for coupling g_a at scales below their mass. To lowest order the couplings scale as

$$\frac{dg_a}{dt} = \frac{b_a}{16\pi^2} g_a^3 \tag{10}$$

$a = 1,2,3,G$ and t is the log of the renormalization point μ. Inputting α and α_s at low energies allows a prediction for M_G, and the weak mixing angle at the weak scale, $\sin^2\theta(M_W)$, in terms of the differences of β function coefficients, $b_a - b_b$. The GUT coupling $\alpha_G = \alpha_G(b_a)$ is also predicted. From Figure 3 we can see that the contributions of the light superpartners have very important effects. Firstly, all couplings are less asymptotically free, so α_G is increased. Furthermore, since α_2 and α_1 run slowly, and they have to grow to quite a large final value, more decades of scaling are required to get there; that is, M_G is increased. Finally, the splitting $b_a - b_b$ are modified in such a way that α_1 and α_2 diverge from each other more slowly, scaling toward the infrared; hence, the weak mixing angle increases.

(b) Yukawa couplings $\lambda_{U,D,E}$

QCD corrections enhance quark couplings to those of the leptons. This gives a successful relation for m_b/m_τ in both ordinary and supersymmetric GUTS. However, this is not particularly compelling since we are absolutely sure there is something wrong or incomplete about the minimal $SU(5)$ Yukawa couplings; they predict $m_e/m_\mu = m_d/m_s$, which is off by a clear order of magnitude.

(c) Gaugino mass parameters

These scale identically to gauge coupling constants, so that $\alpha_a/m_{\tilde{a}}$ is independent of renormalization point μ. This strongly suggests that the gluino is heavier than the photino, probably $m_{\tilde{g}}/m_{\tilde{\gamma}} \simeq 7$.

(d) Scalar mass parameters

The generic equation for the evolution of a squark or slepton mass parameter is schematically

$$\frac{d}{dt} m_i^2 = -a\alpha_a m_{\tilde{a}}^2 + b\,\lambda^2 m_i^2 \tag{11}$$

where a and b are numbers. This is schematic because the flavor structure of the $\lambda_{U,D,E}$ matrices has been ignored. The important point is that large gaugino masses tend to increase m_i^2 in the infrared, so we expect squarks to be heavier than sleptons by an amount which depends on $\tilde{m}$. On the other hand a large Yukawa coupling of i, λ, tends to decrease m_i^2 in the infrared. For three generations, the most important

Yukawa coupling is that of the top. This tends to decrease $m_t^2, m_{t^c}^2$ and m_H^2. A most spectacular result is that, of these, m_H^2 is driven smaller the most rapidly. When it reaches a point at which it is negative it then drives weak symmetry breaking: SUSY breaking is the origin of weak symmetry breaking [17]. While this does not solve all problems, it does make you stop and think: maybe this is even right!

This concludes our brief survey of RG scaling. The full RG equations for the MLES model were written down in reference 18, and a very useful approximate analytic solution appears in the appendix of reference 19. We will return to these RG equations when we discuss flavor.

The RG scaling which we have outlined does suggest various prejudices for the all important question of the spectrum of the light superpartners. There are still many possibilities. One "typical" spectrum is illustrated in Figure 4. It is meant for illustration and may be useful for orientation purposes. It shows the mean squark mass larger than the gluino mass and the photino as the LSP. Neither of these need occur. However the splitting between squarks and sleptons, and the fine structure of these scalar masses due to the D terms, is quite generic. The photino may mix with the other four neutralinos, which are not shown.

II.5 Proton Decay

In both ordinary and supersymmetric grand unified theories proton decay can occur via the exchange of superheavy gauge bosons and scalars (X, Y and H_3, H_3' in $SU(5)$). These yield local four fermion operators so that the proton decay rate is given by

$$\Gamma_p = C\,\alpha_G^2 \frac{m_p^5}{M_{SH}^4} \tag{12}$$

where m_p is the proton mass and M_{SH} the mass of the exchanged superheavy boson. The coefficient C is 0(1), but is uncertain by an order of magnitude [1]. In the minimal $SU(5)$ theory there is an approximate relationship between the superheavy gauge boson mass M_X and the QCD parameter Λ: $M_X \simeq 1.6 \cdot 10^{15}\Lambda$. Even allowing for uncertainties in C, the predicted Γ_p is too large and this theory is excluded.

In passing to supersymmetric $SU(5)$, although α_G is increased, the dominant feature is that M_X is increased by a factor 0(30). Thus gauge boson mediated proton decay is no longer a disaster. On the other hand it is too small to be interesting. The superheavy Higgs triplet exchange is similarly very small if the triplet really is superheavy. It would become interesting if the triplet mass were $0(10^{11}$ GeV).

The most interesting thing about proton decay in supersymmetric theories be-

comes apparent when we perform a general operator analysis [20]. Renormalizable proton decay must be suppressed, as we discussed in the section on R parity. Unlike ordinary unified theories, supersymmetric models have baryon number violating operators with dimension 5: $[QQQL]_F, [U^cU^cD^cE^c]_F$ and three others which violate R. The relevant diagram generating such operators in supersymmetric $SU(5)$ is shown in Figure 5. You might guess that this would lead to the disastrous decay rate $\Gamma_p \simeq \lambda^4 m_p^3/M_{H_3}^2$. This is incorrect because two of the external lines in Figure 5 are scalars and they must be "dressed" at the weak scale to give four fermion operators. A better naive estimate is

$$\Gamma_p \simeq \left(\frac{\lambda^4 m_p^3}{M_{H_3}^2}\right)\left(\frac{\alpha^2 m_p^2}{M_{\tilde{q}}^2}\right). \tag{13}$$

There is enormous literature on a more detailed calculation of this rate, and of all the possible branching ratios. However, I do not believe it can be correct. This decay scenario is inherently coupled to the Yukawa structure of the theory (hence the λ^4 that appears in (13)). The literature that I'm talking of assumes a minimal Yukawa structure ($\lambda_U T T H + \lambda_D T \bar{F} \bar{H}$) which we know to be inadequate from the viewpoint of fermion masses. It does not seem reasonable that it would be adequate to describe proton decay by $d = 5$ operators. This then gives two possible modes of attack on the problem. One is to choose a more complicated Yukawa structure that does give satisfactory fermion masses and mixing angles, and then work out its proton decay rate and modes. The problem here is that there are many such complicated Yukawa structures; but at least it is consistent. An interesting example of this is given in reference 21.

The second strategy is to try to isolate some model independent features of the $d = 5$ proton decay mechanism. The Yukawa nature of the vertices suggests that the flavor structure of the principal decay modes will involve the highest generation kinematically allowed: for example $p \to K^+\bar{\nu}_\mu$ or $K^0\mu^+$. However, this is not a quantitative prediction, and these modes also dominate in certain non-supersymmetric theories with scalars mediated proton decay. I conclude that while supersymmetric $SU(5)$ makes a successful prediction for the weak mixing angle, there is no solid test of this model which we can envisage from proton decay.

II.6 Flavor Physics

You will not need this section for the rest of these lectures. However, it is one of the most interesting aspects of supersymmetric models.

In the standard model, and simple grand unified extensions, flavor physics is confined to the dimension 4 Yukawa couplings $\lambda_{U,D,E}$. The supersymmetric versions of these theories have a much richer flavor structure; there are squark mass operators at dimension 2 and scalar trilinear at dimension 3. In general these can have arbitrary flavor matrix structures: $m^2_{\tilde{Q}}, \ldots$ and $\zeta_{U,D,E}$ for the trilinears. The supergravity boundary conditions imply that these matrices are actually flavor diagonal at the Planck scale. Non-trivial flavor behaviour will then be induced by RG scaling. For example[22]

$$\frac{d}{dt} m^2_{\tilde{Q}} = \frac{1}{8\pi^2} \left\{ \lambda^\dagger_U m^2_{\tilde{U}} \lambda_U + \lambda^\dagger_D m^2_{\tilde{D}^c} \lambda_D \right\} + \cdots, \tag{14}$$

arises from diagrams of Figure 6. These flavor changing mass parameters can have important consequences. A box diagram for $K^0 - \overline{K}^0$ oscillation with internal squarks and gluinos can dominate ϵ and upset the usual determination of KM parameters.

These flavor violating soft operators can be a unique probe of the structure of the theory at extremely high energies.[23] Suppose the complete theory possesses a Yukawa term $\lambda_i Q_i X_1 X_2$ in the superpotential, where $X_{1,2}$ are superheavy. A diagram like that of Figure 6, but with internal squarks replaced by internal X scalars, will lead to a squark mass term with non-trivial flavor structure:

$$\Delta m^2_{\tilde{Q}\,ij} \simeq \frac{\lambda^*_i \lambda_j}{8\pi^2} m^2 \ln(M_p/M_G). \tag{15}$$

In this way, the presence of superheavy particles is felt at low energies by their generation of soft flavor violating couplings. This even occurs in minimal $SU(5)$ where the superheavy scalar in the loop is $\tilde{H}_3$. This leads to lepton number violation in $d = 2$ operators at low energies and induces, for example

$$\Lambda(\mu \to e\gamma) \propto \left| (m^2_{\tilde{E}^c})_{12} \right|^2. \tag{16}$$

In the minimal $SU(5)$ theory the corresponding branching ratio is $0(10^{-15})$ and too small to observe. In more complicated theories it could be much larger.

The much richer flavor structure of supersymmetric theories suggests that they must contain interesting new possibilities for understanding quark and lepton masses and mixing. Although such possibilities exist,[23),24] they have not yet born fruit and so I will not discuss them further.

III. SUPERSYMMETRY AND COSMOLOGY

As mentioned in the introduction, these lectures are directed towards two radical

ideas: that baryogenesis can occur at the weak scale, and that gravitino decay during the keV era can lead to nucleosynthesis compatible with a universe critically dominated by baryons. Both of these ideas have to do with cosmology, and hence in this chapter we give an introduction to elements of the hot big bang cosmology and to conventional supersymmetric cosmology. More complete descriptions can be found elsewhere.[25),26]

In this chapter my equations are approximate. For the sake of clarity I have ignored many subtleties which can be found in the cited references.

III.1 Relics from the Hot Big Bang

The 3^0K cosmic microwave background radiation provides evidence that the early universe was highly isotropic and homogeneous.[27] The expansion of the hot plasma at temperature T is described by the Robertson-Walker metric with scale factor $R(t)$. The expansion rate $\dot{R}/R \simeq \sqrt{\rho}/M_p$ follows from Einstein's equations, with M_p the Planck mass and ρ the total energy density. A relativistic species X in thermal equilibrium with the hot plasma has a number density $n_x(T) \simeq T^3$ so that $\rho_x(T) \simeq T^4$ and $\Gamma_{EXP} \simeq T^2/M_p$ in an era where ρ is dominated by radiation.

A critical question in the hot big bang model is that of relic abundances of heavy, long lived particles. The abundance n_x obeys an equation

$$\dot{n}_x = -\frac{3\dot{R}}{R} n_x - \left(n_x^2 - n_{x_e}^2 \right) \sigma v \tag{17}$$

where n_{x_e} is the equilibrium abundance for X and σv is the thermally averaged annihilation cross-section times relative velocity. Defining a number density normalized to that of photons, $f_x = n_x/n_\gamma$, one has that for $T \ll m_x$: $f_{x_e} \propto e^{-m_x/T}$ as shown in Figure 7. The solution to Equation (17) has the property that f_x tracks f_{x_e} until some point at which a freezeout behaviour occurs.[28] The larger σv is, the lower is f_f, the freezeout abundance. A very simple analytic estimate for this behaviour goes as follows. Freezeout occurs at a temperature T_f at which $\Gamma_{ann}(T_f) = \Gamma_{EXP}(T_f)$, *i.e.* $n_{x_f} \sigma v \simeq T_f^2/M_p$. At lower temperatures the particle number densities are so low that the annihilation mean free time is much longer than the age of the universe. Dividing the above result by T_f^3 gives

$$f_{x_f} = \frac{n_{x_f}}{T_f^3} \simeq \frac{1}{M_p T_f \sigma v}. \tag{18}$$

This is a very useful result. To use it you need to know T_f as well as σv. The freezeout

temperature only depends logarithmically on σv, and for most applications putting $T_f \simeq m_x/20$ is a reasonable approximation.

Consider the case that the $\tilde{e}$ is the lightest superpartner (LSP). In this case the annihilation occurs to two photons with $\sigma v \simeq \alpha^2/m_{\tilde{e}}^2$. This gives $f_{x_f} \sim 0(10^{-12})$ for $m_{\tilde{e}} = 100$ GeV. Since the observable matter in the universe gives a baryon abundance $f_B \simeq 3 \cdot 10^{-10}$ we would have to conclude that more than 1/10 % of all atoms involve $\tilde{e}$ or $\tilde{e}^*$. This seems to be excluded. It is safer to have the LSP be more weakly interacting so that it is unable to dissipate its energy during star formation, and it remains in galactic halos rather than in stars and planets.

If the LSP is a 5 GeV $\tilde{\gamma}$ it will annihilate via t channel scalar exchange to $c\bar{c}$ or $\tau\bar{\tau}$ pairs with $\sigma v \simeq \frac{\alpha^2(3m_f^2+m_{\tilde{f}}^2)}{m^4}$.[29] For values of the scalar mass $m = 100$ GeV, this gives $f_{\tilde{\gamma}} \simeq 10^{-8}$. Since $f_B \simeq 3 \cdot 10^{-10}$, $m_B \simeq 1$ GeV leads to $\Omega_B \simeq .01$, we conclude that $f_{\tilde{\gamma}} \simeq 10^{-8}$, $m_{\tilde{\gamma}} \simeq 5$ GeV leads to $\Omega_{\tilde{\gamma}} \simeq 1$, so that the $\tilde{\gamma}$ is an excellent candidate for the dark matter of the universe.

III.2 The Gravitino Problem

There is a very general conclusion to be drawn from Equation (18). Whenever the annihilation cross section σv for a long-lived particle, at energy scales of order its mass, is extremely small, there will be a cosmological problem because so many of these particles will survive annihilation. In supersymmetry the superpartner of the graviton, the gravitino $\tilde{G}$, is such a particle.[30]

With supersymmetry unbroken the gravitino is massless. In the models which we are considering, those where the supersymmetry breaking operators described in II occur at tree level, the gravitino mass sets the scale for all the supersymmetry breaking and hence is of order the weak scale. The gravitino couples to all particle-superpartner pairs with a coupling that is proportional to momentum p, and the gravitational coupling $1/M_p$. We are interested in annihilations when the momentum is typically less than the gravitino mass, so these couplings are all $0(10^{-17})$. Diagrams such as those in Figure 8 lead to $\sigma v \sim 0(1/M_p^2)$. This is so small that the gravitinos do not annihilate at all. Indeed, they thermally decouple from the hot plasma once the temperature drops significantly beneath M_p. Hence we expect $f_{\tilde{G}} \sim 0(1)$. In fact this is too naive; $f_{\tilde{G}} \sim 0(10^{-2})$ is a better estimate since gravitinos will be diluted by annihilations of all other species.

Two possibilities now arise: $\tilde{G}$ is the LSP and is therefore stable (we are not

considering R_p violation in this section). $\Omega_{\tilde{G}} < 0(1)$ would require $m_{\tilde{G}} < 0$ (keV) which is not the case. The second possibility is that the gravitino is not the LSP and it decays: $\tilde{G} \to q\tilde{q}, g\tilde{g}, \ldots$ Its decay rate $\Gamma_{\tilde{G}} \simeq m_{\tilde{G}}^3/M_p^2$ so that $\tau_{\tilde{G}} \simeq 10^8$ sec $(100\ \text{GeV}/m_{\tilde{G}})^3$. There are many problems with having such an enormous number of heavy particles decaying so late ($T \simeq 100$ eV). At the era of nucleosynthesis the universe may well be matter dominated. This changes the expansion rate and disastrously alters the abundances of the light elements produced. When the $\tilde{G}$ actually decay they produce photons: *e.g.* $\tilde{G} \to \pi^0 \ldots, \pi^0 \to \gamma\gamma$, which do not thermalize and which distort the cosmic background radiation. So many $\tilde{G}$ are decaying that the decay products will lead to destruction of the light elements: *e.g.* $\tilde{G} \to p \ldots, p\ {}^4He \to \ldots$

This is the cosmological gravitino problem. There are two principle ways out: abandon $N = 1$ supergravity models with tree level supersymmetry breaking. In particular if the supersymmetry breaking parameters of Section II arise radiatively, the gravitino mass is no longer linked to the weak scale. Secondly one can use inflation to remove the gravitinos; I explore this second possibility.

Inflation solves the flatness and horizon problems of the universe by having an era in which $R(t) \propto e^{Ht}$ expands exponentially by at least 28 orders of magnitude. This depletes the gravitinos, along with all other particles, to a negligible level, so that immediately after inflation the universe is dominated by inflatons. If these inflatons couple to quarks and leptons they decay at a rate Γ_I and reheat the universe at time $t_R \sim 1/\Gamma_I \sim 1/\Gamma_{EXP} \sim M_p/\sqrt{\rho_I}$ where $\rho_I(t_R) \simeq T_R^4$ where T_R is the temperature to which the universe reheats after inflation. Hence $T_R \simeq \sqrt{M_p\Gamma_I}$ can be chosen to be anything we wish by choice of Γ_I.

Does the process of reheating make any gravitinos? Their couplings are suppressed by M_p^{-1} and hence we do not expect to make many in direct inflaton decays. However, some can be made by scattering processes, as shown in Figure 9. This gives a rate for $\tilde{G}$ production $\Gamma_{\tilde{G}} \simeq \alpha_s T_R^3/M_p^2$. We are assuming that T_R is much larger than any of the particle masses. The number of $\tilde{G}$ made by this process at temperature T_R depends on how much time we have. This time is of order the age of the universe at T_R: $t \sim \Gamma_{EXP}^{-1} \sim M_p/T_R^2$, hence $f_{\tilde{G}} = n_{\tilde{G}}/T^3 \sim \Gamma_{\tilde{G}} t \sim \alpha_s T_R/M_p$.

Clearly we must not have too large a T_R. How large can T_R be? Here is a simplified description of how a limit can be obtained from big bang nucleosynthesis.[31] Suppose $m_{\tilde{G}} \simeq 100$ GeV; in this case its hadronic branching ratio will be of order 1 and it turns out that reactions such as $p\ {}^4He$ will destroy 0(1) ${}^4He$ per $\tilde{G}$ decay. Standard

nucleosynthesis calculations are uncomfortable if more than 10% of the produced ^4He were destroyed so that $f_{\tilde{G}} \lesssim \frac{1}{10} f_{^4He} \simeq 10^{-2} f_B$. Since $f_B \lesssim 10^{-9}$, this gives $f_G \lesssim 10^{-11}$ and $T_R \lesssim 10^9$ GeV. A more careful analysis gives 10^8 GeV $< T_{R_{max}} < 10^{12}$ GeV for $100 < m_{\tilde{G}} < 1000$ GeV. The cosmological gravitino problem has now been rephrased: how can a baryon asymmetry be created with such a low value for T_R?

There is a large literature on generating a baryon asymmetry from inflaton decay at this intermediate scale.[32] The simplest models do not work, and I find the more complicated versions unattractive. Rather than discuss them in detail, in the next section I will describe a completely different approach; one in which baryogenesis is only successful if $T_R \lesssim 0$(GeV)!

IV. Baryogenesis: Conventional and Exotic

IV.1 Baryon Symmetry v. Baryon Asymmetry

It is obvious that locally, on earth, there is an asymmetry between matter and anti-matter. Is this just a local fluctuation, or does the same asymmetry persist throughout the universe? Observationally we can get somewhere toward answering this question because if different domains of the universe have differing signs of the asymmetry, there will be γ rays emanating from the annihilations occurring at the domain walls. From present observation we can conclude that there is very little antimatter within our own local group of galaxies, but it is quite possible that distant clusters are made of anti-matter.[33]

The hot big bang theory leads us to believe that the entire visible universe is comprised of matter, and that there are no visible domain walls. Let us first understand the prediction for the case of a homogeneous universe with no baryon asymmetry. At temperatures much larger than that of the QCD phase transition the light quarks are in thermal equilibrium with the hot plasma with number densities $n_q(T) = n_{\bar{q}}(T) \simeq n_\gamma(T) \simeq T^3$. As the temperature drops through the QCD phase transition these quarks will assemble into color singlet states: $qqq \to B$, $\bar{q}\bar{q}\bar{q} \to \bar{B}$, $\bar{q}q \to \pi, K, \ldots$ The temperature is less than the baryon mass so that $B\bar{B}$ annihilation will occur. The number densities follow the Maxwell-Boltzmann exponential decline with temperature until freezeout occurs at $T_f \simeq m_B/50 \simeq 20$ MeV. Using the non-relativistic $\bar{B}B$ annihilation cross-section in Equation (18) then gives relic abundances $f_B = f_{\bar{B}} \simeq m_B/M_p \simeq 10^{-19}$. This is to be compared with observation of luminous galaxies $f_{B+\bar{B}} > 10^{-11}$.

If we want the universe to have no net baryon number, we must avoid this high

degree of annihilation by having B separated from $\bar{B}$ at $T \gtrsim 20$ MeV, *i.e.* at times earlier than 10^{-3} seconds. Presumably at any time t the size of the domains of baryons and anti-baryons is at most that of the horizon $\simeq$ ct. Thus at 20 MeV the maximum domain mass is $(3 \cdot 10^7 \text{ s.})^3 (20 \text{ MeV})^4 \simeq 10^{56}$ GeV $\simeq .1 M_\odot$. Even if we could prevent these domains from subsequent annihilation, they are far too small to evolve into galaxies. Hence the prejudice that the hot big bang theory requires a primordial baryon asymmetry.[34]

IV.2 Baryogenesis Requirements and the Standard Scenario

In this section we discuss, omitting any detailed proofs, the requirements which most schemes for baryogenesis must fulfill and give the standard scenario: the out of equilibrium decays of a super-heavy scalar in $SU(5)$, as an example.

1. B violation

This is always a necessity. In $SU(5)$ the scalar color triplet partner H_3 of the usual Higgs doublet has B violating couplings. The Yukawa coupling TTH as well as containing QU^cH_2, contains QQH_3 and $U^cE^cH_3$. Similarly $T\bar{F}H^*$ contains $QD^cH_2^*$, QLH_3^* and $U^cD^cH_3^*$. There is no way of defining a conserved baryon number.

2. Decays to states of different B

At temperatures much above the mass of the superheavy triplet, $T \gg M_{H_3} = M$, the triplets are in thermal equilibrium, $n_{H_3} = n_{H_3^c} \simeq T$. As T drips below M the H_3 annihilate and begin to follow an exponential Boltzmann suppression. Those that are left after freezeout, given by Equation (18), will decay. The tree level couplings given above allow

$$\begin{aligned} H_3 &\to \bar{q}\bar{q} \quad \text{with branching ratio } r \\ H_3 &\to q\ell \quad \text{with branching ratio } 1 - r \\ H_3^* &\to qq \quad \text{with branching ratio } \bar{r} \\ H_3^* &\to \bar{q}\bar{\ell} \quad \text{with branching ratio } (1 - \bar{r}) \end{aligned}$$

The baryon asymmetry generated by these decays is then

$$n_B - n_{\bar{B}} \simeq (n_{H_3})_f (\bar{r} - r). \tag{19}$$

Note that if decays always gave a state with the same baryon number, *e.g.* $H_3 \to \bar{q}\bar{q}$ or $H_3^* \to qq$, then no asymmetry can be generated because we are starting with equal numbers of H_3 and H_3^*.

3. CP violation

Is a rather general requirement for the generation of a baryon asymmetry. In the decays of H_3 it is necessary in order that $r \neq \bar{r}$. This is because $H_3^* \to qq$ is the CP conjugate process to $H_3 \to \bar{q}\bar{q}$, so that CP conservation implies $r = \bar{r}$. In the $SU(5)$ theory CP is violated by the Yukawa coupling matrices $\lambda_{U,D}$. To get a large enough η it is typically necessary to introduce further 5-plets and Yukawa interactions.

4. A physical cut

At tree level the amplitudes for $H_3 \to \bar{q}\bar{q}$ and $H_3^* \to qq$ can be written as $\mathcal{A} = \lambda_0 e^{i\theta}$, $\overline{\mathcal{A}} = \lambda_0 e^{-i\theta}$. The phase θ drops out when we compute the rate, so $r = \bar{r}$. In order to get $r \neq \bar{r}$ the result for the rate must know that CP is violated, *i.e.* it must contain the phase θ coming from some coupling constant. This can only occur if there is a loop diagram with an imaginary Feynman integral. Such an absorptive amplitudes occurs when the loop diagram has a cut; *i.e.* there is a contribution from an on-shell intermediate state.

For $H_3 \to \bar{q}\bar{q}$ in the $SU(5)$ theory this is shown in Figure 10. The amplitudes can now be written as $\mathcal{A} = \lambda_1 + \lambda_2 I$, $\overline{\mathcal{A}} = \lambda_1^* + \lambda_2^* I$ where $\lambda_{1,2}$ are coupling constants and I is the loop integral. Then

$$r - \bar{r} \propto |\mathcal{A}|^2 - \left|\bar{\mathcal{A}}\right|^2 \propto \mathrm{Im}(\lambda_1^* \lambda_2)\,\mathrm{Im}\, I. \tag{20}$$

This shows the requirement for the imaginary piece of the loop integral.

5. Amplitudes with different phases

If all contributions to a given amplitude have the same phase, then the phase will not appear in the rate. In the above example, $r - \bar{r} \neq 0$ requires that λ_1 and λ_2 have a different phase. This is a non-trivial restriction for the minimal $SU(5)$ theory.

6. B violating rescattering

This is really a combination of *2.* and *4.* Suppose that the cut and the final state have the same baryon number. In this case the diagrams which have $\mathrm{Im}\, I \neq 0$ do not know that there are decay modes with differing baryon number, hence they do not lead to $n_B \neq n_{\bar{B}}$. When the rescattering between cut and final state violates B, the subset of diagrams with $\mathrm{Im}\, I \neq 0$ also knows that there are possible final states with differing B. In Figure 11 we show H_3 decaying to final state S_1, with a possible final state S_2 in the cut. The rescattering theorem requires that $B(S_2) \neq B(S_1)$ in order for $n_B \neq n_{\bar{B}}$. If $B(S_2) = B(S_1)$ then any η generated in $H_3 \to S_1$ is cancelled by that generated in $H_3 \to S_2$. In the minimal $SU(5)$ theory the diagrams of Figure 10 do

satisfy the rescattering theorem.

7. Out of thermal equilibrium

It is not so easy to prove that baryogenesis requires a system to be out of thermal equilibrium, but we can make it seem plausible. In the $SU(5)$ theory thermal equilibrium would mean that the reactions $H_3 \leftrightarrow \bar{q}\bar{q}$, $H_3 \leftrightarrow q\ell$ are occurring in both directions. Alternatively the conversions $\bar{q}\bar{q} \leftrightarrow q\ell$ and $qq \leftrightarrow \bar{q}\bar{\ell}$ can occur. Imagine that there is a tendency towards producing a quark excess. This is self regulating since then the reaction rates to reduce this excess are enhanced by the increased density of quarks.

The language of chemical potentials allows us to immediately conclude that baryogenesis cannot occur in thermal equilibrium. In thermal equilibrium a particle-antiparticle excess for species i can always be described by a chemical potential $\mu_i(T)$:

$$f_{i,\bar{i}} \propto \left(e^{\frac{E(p)}{kT} \pm \frac{\mu_i(T)}{kT}} \pm 1 \right)^{-1} \tag{21}$$

where the $\pm$ in the exponent refers to i or $\bar{i}$ and the $\pm$ 1 refers to fermion or boson. Reactions in thermal equilibrium lead to additive relationships amongst chemical potentials: $\bar{q}q \leftrightarrow \gamma \Longrightarrow \mu_q = -\mu_{\bar{q}}$, which we have already imposed. Such relationships imply that the number of independent chemical potentials corresponds to the number of quantum numbers conserved by the reactions which are in thermal equilibrium. A system in thermal equilibrium cannot support a non-zero chemical potential for a quantum number which is violated. If it starts out non-zero it will be washed out. Once it is zero it remains zero.

The minimal $SU(5)$ theory is rather interesting because while B and L are both violated, $B - L$ is conserved and therefore possesses a chemical potential. To see this note that the reactions $qq \leftrightarrow \bar{q}\bar{\ell}$ and $\bar{q}\bar{q} \leftrightarrow q\ell$ give $\mu_\ell = -3\mu_q \equiv -3\mu_{B-L}$. You can check that $\mu_{B-L} \neq 0$ corresponds to non-zero $B - L$: if μ_{B-L} is positive there is an excess of quarks and antileptons. However, we are imagining that the universe starts with no asymmetries so initially $\mu_{B-L} = 0$. Since $B - L$ is conserved, μ_{B-L} must always remain zero.

In the standard $SU(5)$ scenario the idea is that the H_3 are out of thermal equilibrium once the relic abundance is frozen out (actually even before freezeout). This is because with $T \ll M$, the decays $H_3 \to \bar{q}\bar{q}, q\ell$ occur on a time scale Γ^{-1}, but the inverse decays are not occurring at all because the thermal quarks and leptons typically have too small energies to produce a H_3 when they collide. The abundances

f_{H_3} and $f_{H_3^c}$ are shown schematically in Figure 12. This picture only makes sense if the H_3 decay rate is smaller than the expansion rate of the universe when $T = M$. This requires $M \gtrsim \lambda^2 M_p$ where λ is the largest Yukawa coupling of H_3 to light states. Hence this conventional mechanism is restricted to decays of very heavy states.

8. Preventing destruction of η

The final requirement is that η is not washed out once it has been made. From our discussion of chemical potentials it should be clear that we must require that after baryogenesis the reactions of the hot plasma which are in thermal equilibrium must conserve baryon number. Only in this way can a non-zero chemical potential for baryons be maintained. This means that the scattering reactions $qq \leftrightarrow \bar{q}\bar{\ell}$ which have a rate $\Gamma_{\Delta B}$ must be out of thermal equilibrium. That is $\Gamma_{\Delta B} \ll \Gamma_{EXP}(t)$ for all t (T) larger (smaller) than that of the baryogenesis era. We now show that this is automatically satisfied in the standard scenario for all $M < M_p$.

At temperature T, $\Gamma_{\Delta B} = n_q(\sigma v)_{\Delta B} \simeq \lambda^4 T^5/M_p^4$. This falls much more rapidly with T than the expansion rate of the universe $\Gamma_{EXP} \simeq T^2/M_p$. Hence a sufficient condition to prevent washout is

$$\Gamma_{\Delta B}(T_D) \ll \Gamma_{EXP}(T_D) \tag{22}$$

where T_D is the temperature at which the H_3 decay: $T_D^2 = \lambda^2 M M_p$. We find $\Gamma_{\Delta B}(T_D)/\Gamma_{EXP}(T_D) \simeq M/M_p$ which is smaller than unity for all $M < M_p$.

We warn the reader that this simple analytic argument hides the true complexities of a detailed calculation. In practice one must use exponential decays rather than the sudden approximation; deviations from thermal equilibrium are small (the top Yukawa is large) and the $2 \to 2$ scattering can often not be ignored. In fact the $2 \to 2$ scattering may end up being more important than decays in generating η.

IV.3 The "$\Delta B \neq 0$" Model and Weak Scale Baryogenesis

In the previous section I discussed certain conditions which allow baryogenesis to occur (they are not all rigorous requirements) and illustrated how these conditions are satisfied in the standard scenario of decays of superheavy bosons. There is a problem with this scenario: will we ever be able to test it? The same problem exists with the many new baryogenesis schemes which have been invented over the last five years or so. In this section I will present a baryogenesis model which has features that are testable, because the physics responsible for baryogenesis is at the weak scale. Such models are hard to come by, and I have not been able to find a model in which

the numerical value for η can be completely predicted from low energy measurements.

Weak scale, and therefore testable, baryogenesis models are hard to construct for several reasons:

i) there is the obvious problem of proton stability, ^{16}O stability, $n - \bar{n}$ oscillations, etc.

ii) The lower the temperature of the universe the harder it is to get out of thermal equilibrium. This is because $\Gamma_{EXP} \propto T^2$, so that at lower temperatures more time is available for reactions to equilibrate.

iii) By the same reasoning, the lower the temperature at which baryon number violating processes occur the harder it is to prevent η being washed out; at low temperatures there is simply more time for $\Delta B \neq 0$ scatterings to destroy an asymmetry.

iv) The numerical value of η often depends on cosmological parameters which are hard to measure. For example, this might be the expansion rate at a certain era, or in the case discussed below it will be the properties of the inflaton.

In the remainder of this section I show how baryogenesis could occur in the R_p violating "$\Delta B \neq 0$" model discussed in Section II.3.[35] This supersymmetric model has minimal field content and a superpotential $f + f_{\Delta B}$ given in Equations (1) and (7). The symmetry structure of the operator $U_i^c D_j^c D_k^c$ is such that λ_{3ijk} is antisymmetric on $j \leftrightarrow k$.

Although our model has a stable proton, there are $\Delta B = 2$ six-quark operators induced which mediate neutron anti-neutron oscillations and the decay of heavy nuclei such as ^{16}O and Fe. The diagram of Figure 13 is the most dangerous contribution to neutron oscillations. The necessity of the flavor-changing scalar masses, forced by the anti-symmetry property of λ_3, means that this provides a less stringent constraint than from heavy nuclei decay. In heavy nuclei decay, kaons can appear in the final state, so the flavor change is not necessary. We estimate that heavy nuclei will have lifetimes greater than 10^{31} yr. providing

$$(\lambda_3)_{112} \lesssim 10^{-6} \left(\frac{m}{300\ \text{GeV}}\right)^{5/2}. \tag{23}$$

This is not an unnaturally small Yukawa coupling to low generations. Note that higher generation couplings, such as $(\lambda_3)_{332} \equiv \lambda$, can be $0(1)$.

We will use inflation to get out of thermal equilibrium since this is needed to remove the large number of gravitinos in any case. The inflationary phase transition leads to an out of thermal equilibrium situation, and baryogenesis will occur at the end of this phase transition. At the end of the inflationary era the vacuum energy which

was driving inflation becomes inflaton rest mass energy. For long-lived inflatons, this energy density is diluted during a matter dominated era. Let ρ_D be their energy density at the era of inflaton decay. We assume that squarks are produced in the direct decay products of the inflaton. The baryon asymmetry will be generated in squark decays. The reheat temperature was calculated in Chapter III to be $T_R \sim \rho_D^{1/4} \sim \sqrt{\Gamma_I M_p}$. This will be chosen to be less than the squark mass (m). Γ_I is the inflaton decay rate and M_p the Planck mass. Thus $M_I > m > T_R$, where M_I is the inflaton mass. The baryon asymmetry produced, expressed as the ratio of baryon to photon number densities today,[36] is

$$\eta \equiv n_B/n_\gamma \simeq 5R(T_R/M_I)\Delta B, \tag{24}$$

where R is the inflaton decay branching ratio to the squarks of interest and ΔB is the average baryon asymmetry produced by the decay of one such squark anti-squark pair.

After the squarks have decayed and generated a baryon asymmetry, the universe rapidly thermalizes at T_R. In order for scattering and inverse decays not to destroy the baryon asymmetry, we find that T_R must be low; less than a few GeV. Such a low T_R implies a model for inflation different from those currently used. However, there is no conflict between low T_R, which depends on Γ_I, and having sufficient expansion and acceptable density perturbations which depend on other features of the inflaton potential. Furthermore such a low T_R implies that essentially no gravitinos are produced after inflation: the gravitino problem is solved in our scheme.

The squarks made in direct inflaton decays will have momenta of order M_I and hence will be far from thermal equilibrium. They will be destroyed either by decays, annihilations or scatterings. For $T_R \ll M_I$, the particle number densities are low so that decays always dominate. For simplicity we now study the baryon asymmetry generated by the decay of just $\tilde{t}^c$ and $\tilde{b}^c$, as would dominate if $(\lambda_3)_{332}$ is the largest element of λ_3. In this case the dominant contribution to ΔB comes from the diagram shown in Figure 14. This diagram represents the interference between two one loop diagrams for $\tilde{t}^c \to \bar{s}^c\bar{b}^c$. We now discuss how this satisfies the eight conditions for baryogenesis listed in the previous section

i) Baryon number is violated by $(\lambda_3)_{332} \equiv \lambda$.

ii) $\tilde{t}^c$ has decay modes of differing baryon number: $\tilde{t}^c \to \bar{b}^c\bar{s}^c$ and $\tilde{t}^c \to \tilde{g}t^c$. Note that this requires $m(\tilde{t}^c) > m_{\tilde{g}} + m_t$.

iii) CP is violated at the trilinear squark vertex which is proportional to $|A|\,e^{i\phi_A}$,

where the A parameter was defined in Equation (2c). Note that this phase ϕ_A is precisely the same as that which led to a potentially large neutron electric dipole moment [Equation (4)].

iv) If we consider $\tilde{t}^c \to \bar{b}^c\bar{s}^c$, the one loop diagram with the intermediate $\tilde{g}t^c$ has a physical cut.

v) The two one loop amplitudes have different phases. One is real, the other is proportional to $e^{i\phi_A}$.

vi) The $\Delta B \neq 0$ rescattering in Figure 14 is $\bar{b}^c\bar{s}^c \to \tilde{g}t^c$.

vii) Inflation has been used to satisfy the out of thermal equilibrium condition.

viii) η washout is the final challenge, and we discuss it shortly.

From Figure 14 we estimate

$$\Delta B \simeq \frac{\alpha_s}{64\pi^2} A\,\phi_A\,|\lambda|^2\,. \tag{25}$$

In order to use these results to calculate $\eta = n_B/n_\gamma$, we must estimate T_R/M_I [see Eq. (24)]. T_R is restricted because it is necessary to ensure that reactions occurring in the hot plasma do not destroy the baryon asymmetry which has been produced. This is a particular worry for low T_R since the universe is expanding slowly and there is much time for the B violating interactions to damp the asymmetry. Once the squarks have decayed, the universe will be brought to thermal equilibrium at rates or order T_R. The baryon number will not be destroyed during this non-equilibrium era because there is insufficient time. After thermalization, the expansion rate is $\Gamma_{EXP} \sim T_R^2/M_p$. The most dangerous processes for erasing the baryon asymmetry are shown in Fig. 15. If $T_R \gtrsim m$, these processes have a rate $\Gamma_{\Delta B} \sim \alpha_s\,|\lambda|^2\,T_R$ which is certainly much larger than Γ_{EXP} and would destroy η. To suppress this rate we require $T_R \ll m_{\tilde{\gamma}}$. In this case diagram (a) dominates, but has a rate suppressed by $\exp[-(m_{\tilde{\gamma}}/T_R)]$:

$$\Gamma_{\Delta B}(t^c b^c \to \bar{s}^c + \bar{\gamma}) \simeq \alpha\,|\lambda|^2\,(m_{\tilde{\gamma}}^2/m^4) m_b^{3/2} m_t^{3/2} \times \exp[-(m_{\tilde{\gamma}}/T_R)]. \tag{26}$$

Requiring $\Gamma_{\Delta B} \ll \Gamma_{EXP}$ gives $T_R < m_{\tilde{\gamma}}/30$. We have been very rough in estimating the numerical coefficient and terms in $\Gamma_{\Delta B}$. This is because the exponential factor in (9) has an overwhelming effect and $T_R/m_{\tilde{\gamma}} < 1/50$ will certainly be safe. In $N = 1$ supergravity theories, $m_{\tilde{g}} \simeq 7m_{\tilde{\gamma}}$. Furthermore $m > m_{\tilde{g}}$ and $M_I > m$. Hence a reasonable lower bound for T_R/M_I is 10^{-3}.

Making use of (25) and (4) in (24) gives

$$\frac{\eta}{5\cdot 10^{-10}} \simeq \left(\frac{R}{1/3}\right)\left(\frac{T_R/M_I}{10^{-3}}\right)\left(\frac{d_n}{2.5\cdot 10^{-25} e\,cm}\right)\left(\frac{m}{300\text{ GeV}}\right)^2\left|\frac{\lambda}{1/3}\right|^2. \tag{27}$$

R is the average number of $\tilde{b}^c + \tilde{t}^c + \tilde{s}^c$ squarks produced per inflaton decay. It is unlikely to be unity. Even if inflatons decay predominantly to third generation squarks, some of the time we would expect $\tilde{t}$ and $\tilde{b}$. We only need to assume that R is not very small.

Our result for η is very striking. The mechanism works: it is quite reasonable that sufficient η is generated. However, it is inextrically linked to d_n: both are induced by the same phase ϕ_A. Even if we push λ up to 1 and $\tilde{m}$ up to a TeV, d_n would still only be two orders of magnitude below the present bound. The reheat temperature, T_R, would be $\sim$ GeV for $T_R/M_I \sim 10^{-3}$, $m \sim 300$ GeV and $M_I \sim$ TeV. However, there is nothing to prevent a T_R as low as a few MeV; indeed this would be more interesting since d_n would then be very close to the present bound.

Does the mechanism have other testable predictions in addition to d_n and heavy nuclei decay? R parity is broken by the operator $U^cD^cD^c$. This means that the lightest superpartner is unstable and that its decays always violate baryon number. One consequence of this is that our model does not have a candidate exotic stable particle for the dark matter.

The most exciting consequence of the coupling of B and R violation is that it may be possible to observe the baryon number violation of the lightest superpartner decays at current and future accelerators.

The LSP decays are, for example

$$\tilde{\gamma} \to qqq \quad , \quad \overline{qqq}, \tag{28a}$$

$$\tilde{\nu} \to \nu qqq \quad , \quad \nu\overline{qqq}. \tag{28b}$$

There is no longer a cosmological argument that the LSP be color and electromagnetically neutral. Hence other possible LSP decays are

$$\tilde{g} \to qqq \quad , \quad \overline{qqq}, \tag{28c}$$

$$\tilde{e} \to eqqq \quad , \quad e\overline{qqq}, \tag{28d}$$

$$\tilde{q} \to \overline{qq}. \tag{28e}$$

If the LSP is either $\tilde{\nu}$, $\tilde{e}$ or $\tilde{q}$ and has a mass less than $M_Z/2$ there will be very spectacular Z decay events. Note that the usual limits on superpartner masses no longer apply: $e^+e^- \to \gamma\tilde{\gamma}\tilde{\gamma}$ no longer gives a (single photon + missing) signature

and squark pair production at hadron colliders no longer gives monojet events. For example, if the $\tilde{e}$ was the LSP we would get

$$Z \to e^+e^- + \text{ jets}. \tag{29}$$

The jets here are very spherical since they come from b primary quarks. The e^+e^- invariant mass would typically be 20-25 GeV and they may not be very isolated. Hence these events would have been missed at the CERN collider even if they contribute $\sim$ 1% to the Z width. On the other hand they would be unmistakeable at SLC or LEP.

As a second example suppose the $\tilde{\gamma}$ is the LSP. In this case $e^+e^- \to \tilde{\gamma}\tilde{\gamma}$ would give interesting spherical jet events above threshold, and if $m_{\tilde{\gamma}} < M_Z/2$ there is a rare Z decay $Z \to \tilde{\gamma}\tilde{\gamma}$ which could have a branching ratio of order 10^{-5}.

In all these decays it is important to stress the unusual feature of the jets. They come from primary quarks from high generation, and they also violate baryon number! Thus the events of Equation (29) may well have 6 units of strangeness and two units of baryon number. The baryon number violation could be detected statistically by the large number of events with $\Lambda\Lambda$ or $\overline{\Lambda\Lambda}$. This is shown schematically in Figure 16.

In this chapter I have discussed baryogenesis. The idea that a cosmic baryon asymmetry can be understood in the hot big bang cosmology with gauge theory models has been an exciting development of the last decade. There are now so many scenarios for baryogenesis that the model builder risks speculating with untestable ideas. I have tried to show that this need not be the case. Baryogenesis (and inflation!) may occur at the weak scale. This fits quite naturally into an extension of the standard supersymmetric model which has R_p explicitly broken. These models can be tested, they lead to spectacular signatures at present machines if the LSP is light enough, and at planned machines otherwise.

V. $\Omega_B = 1$ AND THE GRAVITINO

Ω is the ratio of the mean density of the universe today to that which would precisely close the universe, and Ω_B is the contribution to Ω from baryonic material. The title of this chapter should be a complete mystery; what on earth could the gravitino have to do with Ω_B being unity? Furthermore the connection involves both nucleosynthesis of the light elements and R_p violation. Connections such as these are very speculative, but they make for very exciting physics. Giving a summary of the ideas now would just give the game away, so I will just let the story unfold in the following sections.

V.1 Nucleosynthesis in the Hot Big Bang

Nucleosynthesis is one of the cornerstones of the hot big bang cosmology. It allows the prediction of four nuclide abundances (^2H, ^3He, ^4He, ^7Li) in terms of essentially one parameter: η. In this section I will give a very brief summary of the main features of standard big bang nucleosynthesis. For further details see references 25, 26, 27, 37 and references therein.

At a temperature of a few MeV the universe contains $e^{\pm}$, ν, $\bar{\nu}$, γ and the excess p and n. These are all in thermal equilibrium: $e^{\pm}$ and p via γ exchange, and ν, $\bar{\nu}$ and n via W and Z exchange. The neutron to proton ratio is therefore given by $n/p = \exp(-(m_n - m_p)/T)$. At a temperature T_f close to 1 MeV the weak interactions freezeout; the neutrinos and neutrons decouple from the now purely electromagnetic plasma. The n/p ratio is frozen at the above value with $T = T_f$.

Beneath an MeV several processes occur. As the temperature drops beneath m_e the $e^{\pm}$ density is reduced by $e^+e^- \rightarrow \gamma\gamma$. The n/p ratio decreases due to neutron decay $n \rightarrow p\bar{\nu}e$. This is a small effect at $T = 1$ MeV ($t = 1$ sec. compared with $\tau_n \simeq$ 1000 seconds), but is substantial by $T = .1$ MeV ($t = 100$ seconds). Also deuterium is frequently formed by $pn \rightarrow d\gamma$. However the buildup of d is quite limited because the reverse reaction $\gamma d \rightarrow pn$ has the benefit of 10^9 photons per baryon so that it is competitive even when the γ are far out in the Boltzmann tail. Once $T < .1$ MeV the Boltzmann tail is too small to prevent substantial d production and further nuclear processing then occurs rapidly through such reactions as: $dn \rightarrow {}^3H\gamma$, $dp \rightarrow {}^3He\gamma$, $n\,{}^3He \rightarrow {}^4He\gamma$ and $^4He\,{}^3He \rightarrow {}^7Be\,\gamma$ followed by $^7Be\,\beta$ decay to ^7Li. Once T drops below .03 MeV these reactions cease: all neutrons have been processed into charged nuclei, and Coulomb barriers prevent further reaction.

The final abundances are sketched in Figure 17. The numbers are from the numerical analysis of Yang *et al.*[38] Observation appears to prefer η in the region shown. We can easily understand the large η behaviour. Increased baryon densities allows the d bottleneck to be overcome more easily since $np \rightarrow d\gamma$ is increased relative to the back-reaction. Hence the onset of nucleosynthesis is earlier, there are more n since fewer have decayed, and the final ^4He abundance increases as this is where the majority of n end up. Since the reaction rates increase the processing of d and ^3He is more complete via $dp \rightarrow {}^3He\gamma$, $n\,{}^3He \rightarrow {}^4He\gamma$; hence ^3He and especially d abundances are reduced. Finally the ^7Li abundance is increased because reactions such as $^4He\,{}^3H \rightarrow {}^7Li\,\gamma$ are able to occur earlier when they are less inhibited by the Coulomb barrier. For a universe with $\Omega_B = 1$ we must be able to

i) Deplete ^4He by 10%.
ii) Deplete ^7Li by a factor 0(100).
iii) Find a source of deuterium which gives $d/p \gtrsim 0(10^{-5})$.

This looks quite hopeless and you may justifiably ask "why bother?" After all if we choose $\Omega_B = .03$ everything appears to work very well. The answer has to do with our prejudice from inflation that $\Omega = 1$. If we accept nucleosynthesis arguments for $\Omega_B \lesssim .1$, then we are forced to invent new physics to supply some stable non-baryonic dark matter. This of course is a very exciting possibility. However, we should also explore the possibility that $\Omega_B = 1$, the dark matter is baryonic and the new physics to be invented is that which readjusts the light nuclear abundances as detailed in *i)* – *iii)* above. The amazing and surprising thing is that this can be done by introducing a single, heavy, long-lived but unstable particle.[39]

V.2 keV Era Nucleosynthesis

Conventional nucleosynthesis ends when $t \simeq 1000$ seconds. Our idea is to introduce a particle X which decays after this and which results in the following processes: a) X has decays to γ, e, π, K, ... which lead to electromagnetic showers. The photons produced can lead to photo destruction of nuclei. b) X has baryonic decays to p, $\bar{p}$, n, $\bar{n}$, ... which produce a baryonic shower in which nuclei can be hadrodestroyed, *e.g.* $p\,{}^4He \rightarrow {}^3He\,d$.

It should be clear that a) and b) have reactions capable of yielding *i)*, *ii)* and *iii)*. However, the details of how it works are very non-trivial; for example, it is not obvious that ^7Li can be photodestroyed by a factor 100 without affecting ^4He. These details are interesting but lengthy. All we need here are the results.

^4He is depleted by hadrodestruction. For $M_X = 1$ TeV about 10 ^4He are destroyed per baryonic X-decay. This constrains

$$\frac{n_X}{n_B} r_B \simeq 10^{-3} \tag{30}$$

where n_X is the number density of X before they decay and r_B is the branching ratio of X to final states containing baryons.

The abundances for d, ^3He, ^6Li and ^7Li are given by a competition between hadroproduction and photodestruction. This can be seen in the simplified rate equations

$$\dot{n}_i \simeq r_B \zeta_i \dot{n}_X - \zeta_\gamma n_i \sigma_{\gamma_i} \dot{n}_X \tag{31}$$

where ζ_i is the number of species i produced per baryonic X decay and ζ_γ the number of photons produced per X decay capable of photodissociating i with cross-section σ_{γ_i}. For parameters of interest these equations hit fixed points ($\dot{n}_i = 0$) thus erasing the prior primordial values and replacing them by

$$n_i \simeq \frac{r_B \zeta_i}{\zeta_\gamma \sigma_{\gamma_i}}. \tag{32}$$

The largest i dependence is in ζ_i. We can therefore predict abundance ratios for these elements purely from the nuclear physics of the shower evolution (*i.e.* no dependence on X properties) with the result[39]

$$d : {}^3He : {}^6Li : {}^7Li \simeq 1 : 1 : 10^{-5} : 10^{-6}. \tag{33}$$

For d, ^3He, and ^7Li we get agreement with observation of these ratios. This is very non-trivial. To obtain correct absolute magnitudes we must make a choice of $r_B/\zeta_\gamma \propto r_B/M_X$:

$$\frac{1}{200} < r_B \frac{\text{TeV}}{M_X} < \frac{1}{10}. \tag{34}$$

Note that we have a prediction that primordial ^6Li production is larger than that of ^7Li. This is not excluded. Although the observed lithium is ^7Li, it is observed in environments where the more fragile ^6Li would have been destroyed. There is one further constraint on the X particle parameters for our scheme to work:

$$2 \cdot 10^5 \text{ sec.} \ < \tau_X < 9 \cdot 10^5 \text{ sec.} \tag{35}$$

This ensures that photons in the electromagnetic shower can destroy d, ^3He, ^6Li and ^7Li, but are not energetic enough to destroy ^4He.

We have obtained $\Omega_B = 1$ by introducing a particle X with properties n_X, r_B, M_X, τ_X constrained by three equations (30), (34) and (35). We can draw three immediate conclusions:

1) Since $r_B \leq 1$, $M_X < 200$ TeV and this new particle cannot be pushed to an arbitrary high energy scale. In fact masses near a TeV are much more likely than this very conservative limit.

2) Since $r_B \leq 1$, $n_X/n_B \gtrsim 10^{-3}$ so that a considerable number of these particles must survive the hot big bang.

3) τ_X of a few 10^5 seconds is very long for such a heavy particle. This suggests that the physics responsible for the decay is either a dimension 6 operator from the intermediate ($M_I \simeq 10^{10}$ GeV) scale: $\Gamma \simeq M_X^5/M_I^4$, or a dimension 5 operator from

the GUT or Planck scale: $\Gamma \simeq M_X^3/M_p^2$. Now, at last, you can understand the title to this chapter: we are going to identify X as the gravitino. Instead of having a cosmological gravitino problem we will use the gravitino to make nucleosynthesis consistent with $\Omega_B = 1$.

V.3 Nucleosynthesis from Gravitino Decays

In this final section we want to answer the following questions. Can the gravitino satisfy the conditions (30), (34) and (35)? If so do we learn any interesting constraints on the parameters of the underlying supersymmetric model. This has not yet been worked out in the literature. Below I give you my preliminary thoughts; you may come across some interesting alternatives yourself by molding a supersymmetric model to fit the constraints.

As discussed in Chapter III, the gravitino decays to all kinematically accessible particle-superpartner pairs, $\gamma\tilde{\gamma}, g\tilde{g}, q\tilde{q}\ldots$ Its lifetime is

$$\tau_{\tilde{G}} \simeq 3 \cdot 10^5 \text{ seconds} \left(\frac{500 \text{ GeV}}{m_{\tilde{G}}}\right)^3 \frac{1}{n} \tag{36}$$

where n is the effective number of open channels. If $\tilde{G} \to g\tilde{g}$ or $q\tilde{q}$ is open, then r_B will not be far from unity. In this case (34) $\Rightarrow m_{\tilde{G}} > 10$ TeV so that $\tau_{\tilde{G}}$ cannot satisfy (35). The gravitino must be lighter than squarks and gluinos. Hence we expect some, or all, of the decays $\tilde{G} \to \gamma\tilde{\gamma}, e\tilde{e}, \nu\tilde{\nu}, \ldots$ In this case n is not a very large number so that, to obtain the desired $\tau_{\tilde{G}}$, $m_{\tilde{G}}$ will be within a factor of 2 or so of 500 GeV. Hence we must arrange for r_B to be within a factor of 2 of the range $1/400 < r_B < 1/20$. Note that the abundance of gravitinos, given by Equation (30), can always be satisfied by choosing the reheat temperature after inflation. We will need T_R of about 10^{11} GeV (apparently the ideas of Chapters 4 and 5 are not to be combined!).

Let us try to get $r_B \sim 0(10^{-2})$ in the MLES model with unbroken R_p. Clearly $\tilde{G} \to \nu\bar{\nu}$ is irrelevant for our purposes; it gives neither electromagnetic nor hadronic showers. $\tilde{G} \to \tilde{\gamma}\gamma, \tilde{e}e$ could provide the dominant electromagnetic decays: $\tilde{G} \to \tilde{\gamma}\gamma^*, \gamma^* \to \bar{q}q$ where γ^* is an off-shell intermediate photon, so the decay is really three body: $\tilde{G} \to \tilde{\gamma}\bar{q}q$. This branching ratio looks too small, $0(10^{-4})$, but it might work if all the numbers were pushed. It seems safer to stick with 2 body decays.

Consider the decay to a neutralino, which is a $(\tilde{\gamma}, \tilde{Z}, \tilde{H})$ combination and its partner. Some fraction of the time the partner is Z or H which we expect to have predominantly hadronic decays. We will take the neutralino to be the LSP since

otherwise it will decay, typically hadronically, to give $r_B \simeq 1$. Clearly we would like this LSP to be mainly $\tilde{\gamma}$, but to have 10% of $\tilde{Z}$ or $\tilde{H}$ mixed into it. To see whether this is reasonable we must study the neutralino mass matrix:

$$\begin{pmatrix} 0 & \mu & 0(M_Z) & \\ \mu & 0 & & \\ 0(M_Z) & & m_{\tilde{2}} & 0 \\ & & 0 & m_{\tilde{1}} \end{pmatrix}$$

which we have shown in the $\tilde{H}$, $\tilde{H}'$, $\tilde{w}_3$, $\tilde{b}$ basis. Various possibilities are obvious: if $\mu \gg M_Z$ the LSP is basically $\tilde{b} = \cos\theta\tilde{\gamma} + \sin\theta\tilde{Z}$. This will give $r_B \simeq \sin^2\theta \simeq .22$ which is too large. If we try $\mu \ll M_Z$ the LSP is mainly $\tilde{H}$ and this certainly does not work. If we make $m_{\tilde{1}}$, $m_{\tilde{2}} \ll M_Z$ the LSP is dominantly photino. Although this would seem to be what we want, it doesn't really work too well: the photino is now much lighter than the Z and hence the gluino is expected to be lighter than $\tilde{G}$ which will give us $r_B \simeq 1$! Hence we are forced into the difficult case of $\mu \simeq M_Z \simeq m_{\tilde{1}} \simeq m_{\tilde{2}}$. There may well be something that works, but it does not seem hopeful: there is no reason why the LSP should have only small $\tilde{Z}, \tilde{H}$ components.

There is, however, a route which looks much better. That of R_p violation. In particular we can look at the "$\Delta L \neq 0$" model. Suppose for example that only τ number is broken so that

$$f_{\Delta L} = \lambda_1 L_1 E_1^c L_3 + \lambda_1' L_2 E_2^c L_3 + \lambda_2 Q D^c L_3.$$

Furthermore I will take $\tilde{\nu}_\tau$ as the LSP. This is a favored case. The $SU(2) \times U(1)$ D^2 terms of the potential tend to push the sneutrinos lighter than the charged scalars, and of the three sneutrinos the $\tilde{\nu}_\tau$ should be the lightest due to the larger RG scaling effects of the τ Yukawa coupling. Thus we have $\tilde{G} \to \nu\bar{\nu}$ followed by $\tilde{\nu} \to e^+e^-$ or $\mu^+\mu^-$, or $\tilde{\nu} \to \bar{q}q$. To get r_B of 10^{-2} we simply choose λ_2/λ_1 or λ_2/λ_1' to be 10^{-1}.

I find this scheme, of $\tilde{G}$ decays readjusting light nuclear abundances so that $\Omega_B = 1$, to be very attractive. It solves two cosmological problems (gravitino, and dark matter) at once. The gravitino is the perfect X candidate—we expect it to have the desired mass and lifetime. Breaking R_p seems to be the thing to do since otherwise the LSP would be stable and would contribute to Ω so that we would not want $\Omega_B = 1$. It is also the smoking gun: R_p violation can be tested by searching for exotic events at present and future e^+e^- and hadron colliders. For example, $\tilde{\nu}_\tau$ would be discovered as

a huge resonance in $e^+e^- \to \tilde{\nu}_\tau \to e^+e^-$.[40] It could give a signal in Bhabha scattering up to a hundred times larger than the Z.

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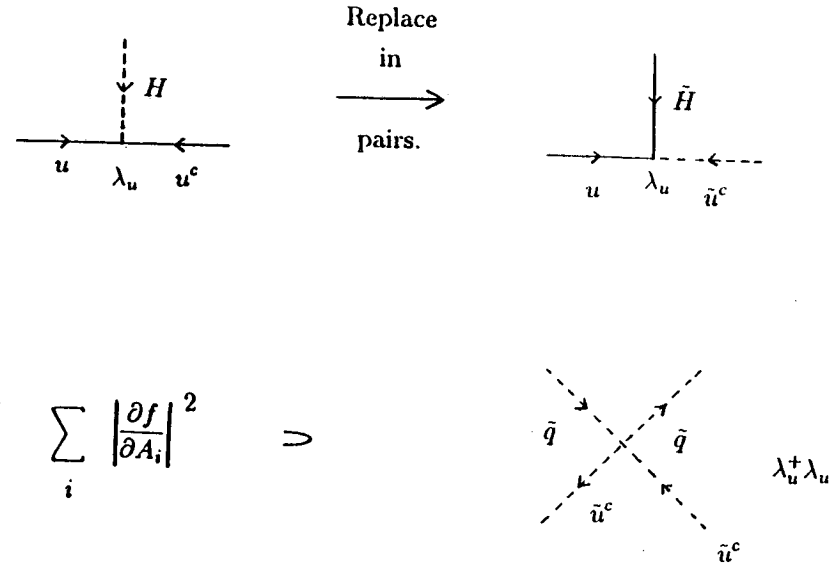


Figure 1 Mnemonics for writing the vertices of the minimal low energy supersymmetric model.

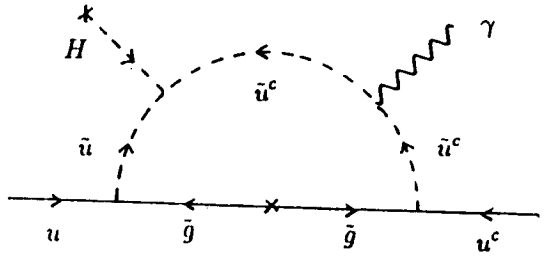


Figure 2 One loop contribution to d_n which is proportional to ϕ_A.

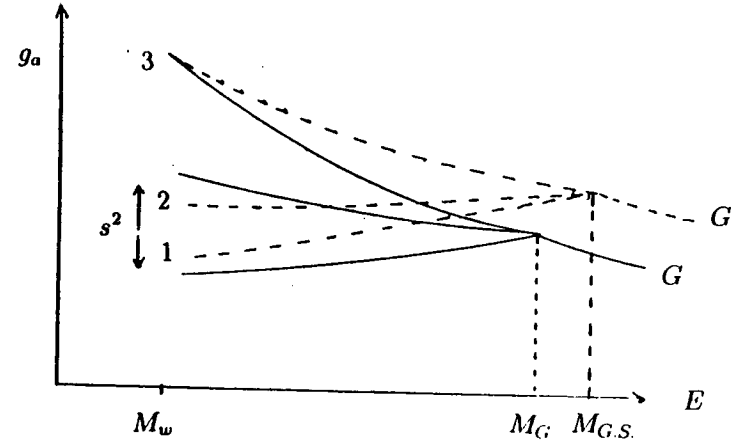


Figure 3 Renormalization group trajectories of gauge coupling constants in ordinary and supersymmetric unification schemes.

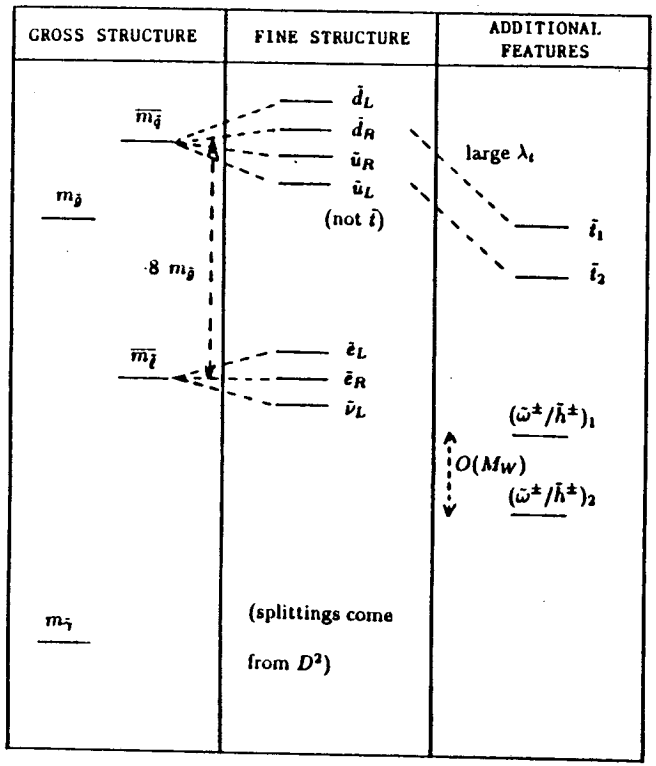


Figure 4 An example of a typical superpartner spectrum ($M_W \ll m, \tilde{m}, \mu$; $\tilde{\gamma}$ as LSP).

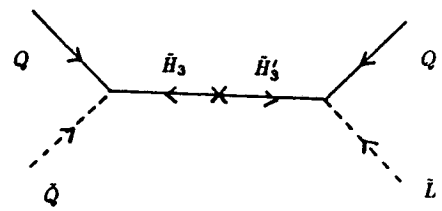


Figure 5 Diagram generating a dimension 5 baryon number violating effective Hamiltonian in supersymmetric $SU(5)$.

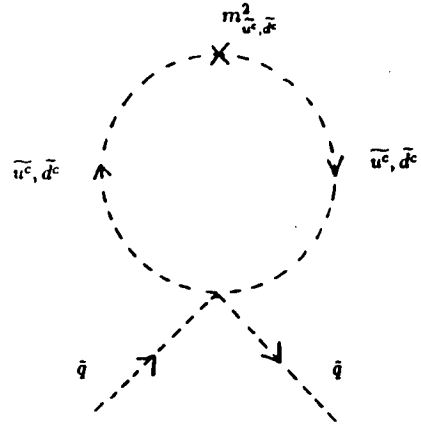


Figure 6 A diagram generating flavor violating scalar masses by renormalizaton group scaling.

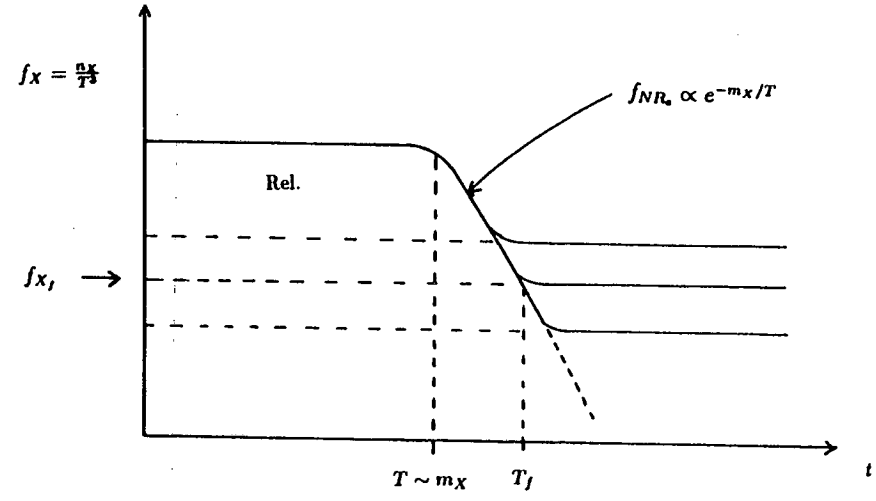


Figure 7 Cosmological freezeout behaviour for the decoupling of a species when it is non-relativistic.

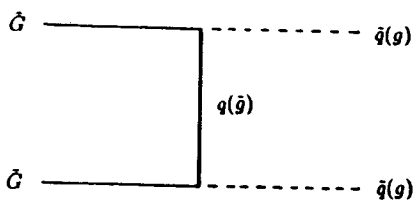


Figure 8 Examples of gravitino-gravitino annihilation processes.

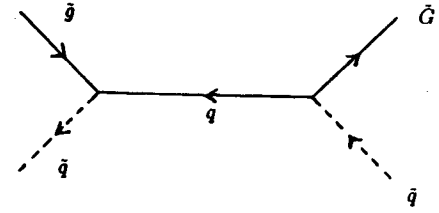


Figure 9 Gravitino production after inflation.

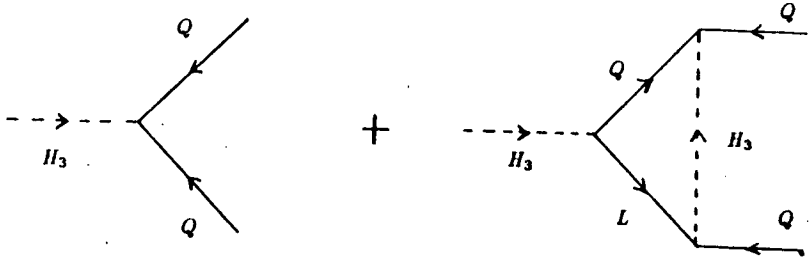


Figure 10 Tree and one loop contributions to the decay of the superheavy scalar triplet $H_3 \to \bar{q}\bar{q}$.

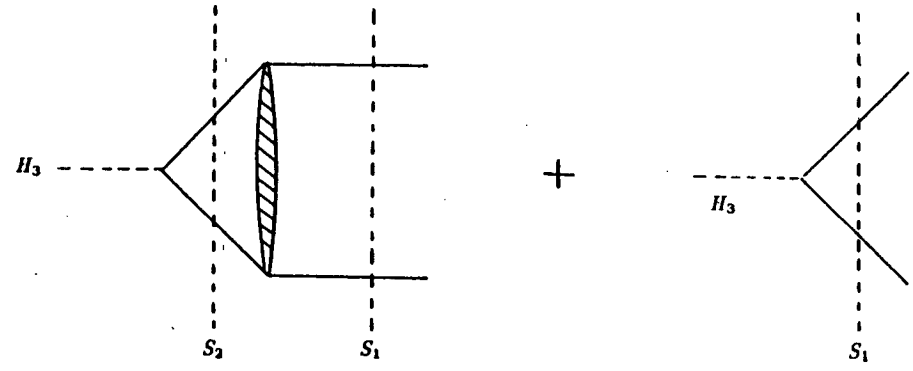


Figure 11 The rescattering theorem: baryogenesis requires $B(S_1) \neq B(S_2)$.

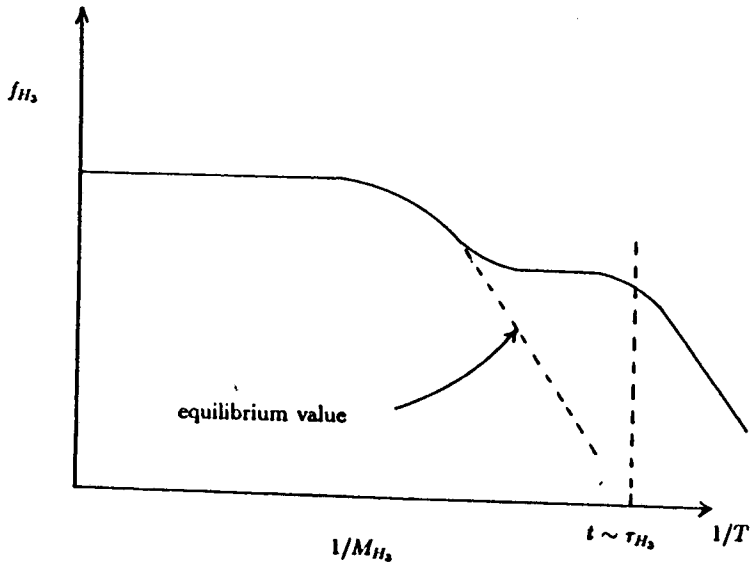


Figure 12 Abundance of superheavy color scalars as a function of temperature. Of those which survive annihilation at $T \simeq M_{H_3}$, most decay out of thermal equilibrium when $t \simeq \tau_{H_3}$.

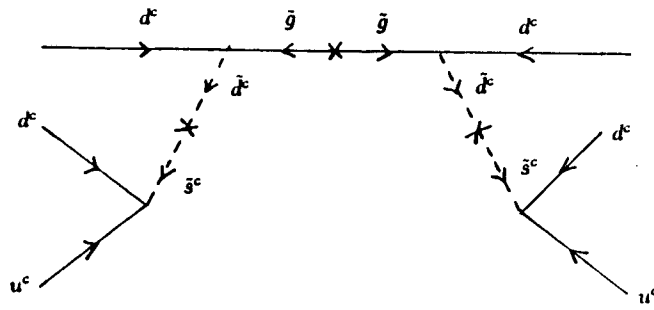


Figure 13 Diagram which induces neutron anti-neutron oscillations.

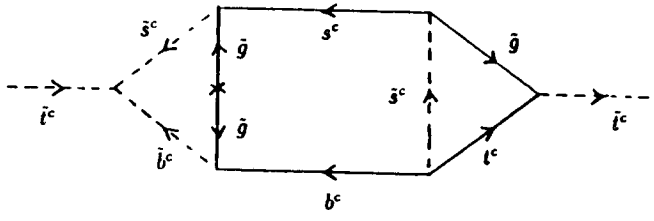
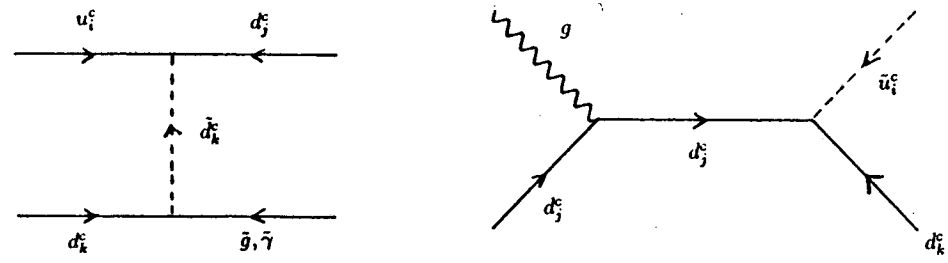


Figure 14 Interference diagram for ΔB in top squark decay.



(a) (b)

Figure 15 Baryon number violating scattering which could lead to the destruction of the baryon asymmetry.

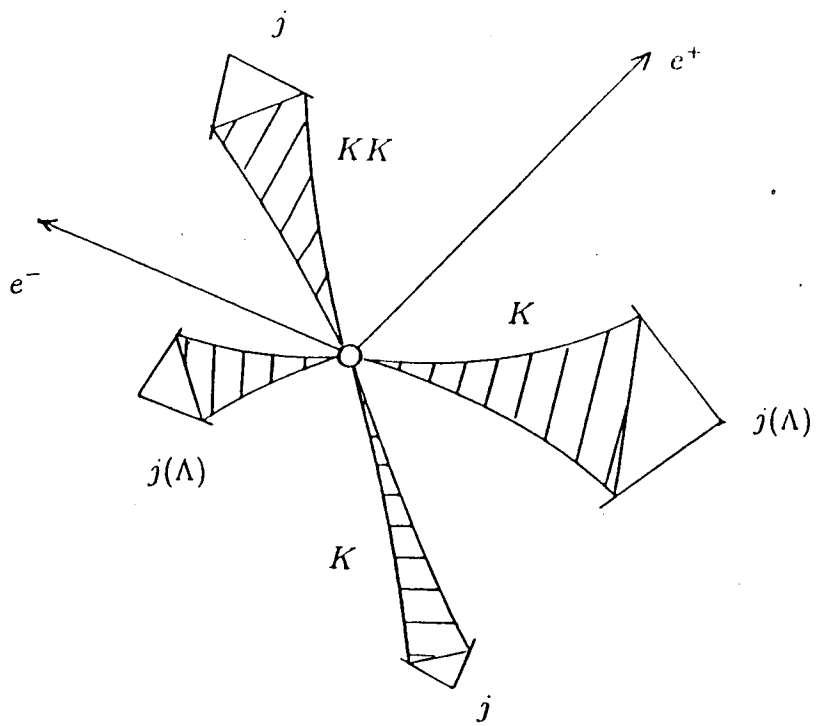


Figure 16 Event topology for $e^+e^- \to Z \to \bar{c}^+\bar{c}^-$. The jets frequently contain $\Lambda\Lambda$, and contain more K than in typical hadronic Z decays.

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