Baryons and Dark Matter from the Late Decay of a Supersymmetric Condensate

SCOTT THOMAS*

Stanford Linear Accelerator Center Stanford University Stanford, CA 94309

Abstract

The possibility that both the baryon asymmetry and dark matter arise from the late decay of a population of supersymmetric particles is considered. If the decay takes place below the LSP freeze out temperature, a nonthermal distribution of LSPs results. With conserved R parity these relic LSPs contribute to the dark matter density. A net asymmetry can exist in the population of decaying particles if it arises from coherent production along a supersymmetric flat direction. The asymmetry is transferred to baryons if the condensate decays through the lowest order nonrenormalizable operators which couple to R odd combinations of standard model particles. This also ensures at least one LSP per decay. The relic baryon to dark matter densities is then naturally $\Omega_b/\Omega_{\rm LSP} \sim \mathcal{O}(m_b/m_{\rm LSP})$. The resulting upper limit on the LSP mass is model dependent but in the range $\mathcal{O}(30-140)$ GeV. The total relic density is related to the order at which the flat direction which gives rise to the condensate is lifted. The observed density is obtained for a direction which is lifted by a fourth order Planck scale suppressed operator in the superpotential.

Submitted to Physics Letters B

Work supported by by Department of Energy contract DE-AC03-76SF00515

INTRODUCTION

The baryon asymmetry and dark matter density may provide indirect windows to very early epochs in the evolution of the universe, and to physics at large energy scales. In most scenarios the physical mechanisms which give rise to the baryon asymmetry and dark matter are unrelated. For example, in supersymmetric theories the dark matter density is usually assumed to result from the freeze out of the lightest supersymmetric particle (LSP). If R parity is unbroken the LSP is stable, and the relic LSPs make up the dark matter. The baryon asymmetry is usually assumed to arise either at the electroweak phase transition [1], by the Affleck-Dine mechanism in which a coherent condensate carrying baryon number is generated along a supersymmetric flat direction [2], or the out of equilibrium decay of massive particles through baryon and CP violating interactions [3]. In all these mechanisms the dark matter and baryon densities are a priori unrelated. This is not surprising since the LSP carries a multiplicative quantum number while baryon number is additive. The processes which lead to the respective relic densities are therefore distinct. That the baryon and dark matter densities are in fact the same within a few orders is not necessarily a direct consequence of any of these mechanisms, and seems fortuitous.

Here I suggest an alternate supersymmetric mechanism in which both baryons and dark matter-arise from the late decay of a weak scale mass particle. As discussed below if the mass of the decaying particle is above the LSP mass, and the population of decaying particles carries a large asymmetry, then (optimally) roughly equal numbers of baryons and LSPs result from the decay. If the temperature at the era of decay is low enough, the LSPs do not rethermalize, and the relic density is determined by the decaying particle density. The ratio of baryon to dark matter density in this scheme is then proportional to the ratio of the lightest baryon mass to LSP mass, $\Omega_b/\Omega_{\rm LSP} \sim \mathcal{O}(m_b/m_{\rm LSP})$. For an LSP with weak scale mass, this gives roughly the correct ratio, $\Omega_b/\Omega_{\rm LSP} \sim \mathcal{O}(10^{-1} - 10^{-2})$. This result is reminiscent of the analogous relation in technicolor theories if the lightest technibaryon makes up the dark matter. There the electroweak anomaly ensures that at high temperatures the baryon and technibaryon number densities are roughly equal [4]. Here however, the LSP density is protected from erasure by the low temperature at the time of decay, rather than an additive quantum number.

In order for this mechanism of relating the baryon and dark matter densities to be operative the decay must occur below the LSP thermalization temperature, but above the temperature at which nucleosynthesis takes place. This can happen if the decaying particle is coupled to standard model fields by nonrenormalizable operators suppressed by an intermediate scale, somewhat below the GUT scale [5]. These operators must carry baryon number if any asymmetry is to result, and be odd under R parity if at least one LSP is to result from each decay. In addition, there should be a large particle-antiparticle asymmetry in the decaying population if order one baryon per decay is to result. Such large asymmetries can result from the coherent production of scalar fields along supersymmetric flat directions. Flat directions are likely to be generic features of supersymmetric theories. Finally, in order that the total density of the universe have the observed value now, the number density of the late decaying particles should be less than thermal at the time of decay. Far too many LSPs would remain if the decaying particles had thermal number density. A subthermal number density in fact naturally occurs for coherent production along flat directions which are lifted by Planck scale suppressed terms in the superpotential [6]. The density in the condensate, and therefore the total relic density, is related to the order at which the flat direction is lifted. All the ingredients for this late decay scenario therefore exist in supersymmetric theories.

REQUIREMENTS FOR BARYONS AND LSPS FROM LATE DECAY

A number of requirements must be met if the late decay scenario for the origin of the baryon asymmetry and dark matter is to be realized within supersymmetry. In most SUSY models the LSP is typically a neutralino, a linear combination of gaugino and Higgsino. If the relic LSPs are to act as dark matter, they must be stable as the result of some symmetry. Since the neutralino is Majorana, this must be a discrete symmetry, giving a multiplicative quantum number. In what follows I will assume the required symmetry is R parity. If the decaying particle is much heavier than the LSP then multiple LSPs can in principle be produced in the decay chain. However, as discussed below at most one unit of baryon number can result from each decay. So unless the LSP is very light, there should not be too many LSPs per decay. In order to guarantee that at least one LSP results from each decay, the decaying particle should be odd under R parity from the low energy point of view. If the mass of the decaying particle is in the range $m_{\text{LSP}} < m_{\phi} < 2m_{\text{LSP}}$, then precisely one LSP results per decay. For simplicity this will be assumed to be the case. The decaying particle then also has weak scale mass.

If the number density of relic LSPs is to be determined by the density of decaying particles, the temperature during the decay epoch should be less than the LSP equilibration temperature. If the decay takes place above this temperature, the relic LSP density is determined by freeze out, as in the usual scenario. For particles with weak scale annihilation cross section and mass, the equilibration temperature is roughly $T \sim \frac{1}{20} m_{\text{LSP}}$. With the LSP mass in the range discussed below this corresponds to roughly $T \sim \mathcal{O}(1 \text{ GeV})$. In addition to this upper limit on the temperature at the time of decay there is a lower limit arising from nucleosynthesis. If decays take place during or after nucleosynthesis the light element abundances can be modified by photodissociation and photoproduction by decay products [7,8]. This can be avoided for $T \gtrsim 1$ MeV since the weak interactions are in equilibrium and the usual neutron to proton ratio results. The decay temperature must therefore lie in the window 1 MeV $\lesssim T \lesssim 1$ GeV. The decay rate, Γ , and decay temperature, T_d are related by $T_d^2 \sim \sqrt{90/g_*\pi^2}\Gamma M_p$, where $M_p = m_p/\sqrt{8\pi}$ is the reduced Planck mass, and g_* is the effective number of degrees of freedom $(g_* \simeq 10.75 \text{ for } T \sim 1 \text{ MeV})$. With weak scale mass, such a slow decay rate implies the decaying particle must couple to standard model fields only through nonrenormalizable interactions. Decay through Planck scale suppressed couplings leads to a decay temperature much too low to avoid the bounds from nucleosynthesis [9,10]. However, a decay temperature of order the nucleosynthesis bound in fact results if the particle decays through dimension 5 operators suppressed by a scale somewhat below the GUT scale [5]. For the 3-body decays discussed below

$$\Gamma \simeq \frac{6\bar{\lambda}^2 m_{\phi}^3}{(8\pi)^3 M^2} \tag{1}$$

where $\bar{\lambda}^2 = \sum |\lambda|^2$ is a sum over generations in the final state, λ/M is the coefficient of the operator, m_{ϕ} is the mass of the decaying particle, and final state masses have been neglected for simplicity. This gives a decay temperature of

$$T_d \sim .3 \left(\frac{10^{14} \text{ GeV}}{M/\bar{\lambda}}\right) \left(\frac{m_{\phi}}{100 \text{ GeV}}\right)^{3/2} \text{ MeV}$$
 (2)

A decay temperature in the window given above can be obtained for 3×10^{10} GeV

 $\lesssim M/\bar{\lambda} \lesssim 3 \times 10^{13}$ GeV. Although this is probably too low to be associated directly with the GUT scale, it could arise from an intermediate scale.

Producing a baryon asymmetry in the decay imposes a number of additional requirements. The particle must of course decay through an operator which transforms under $U(1)_B$ with respect to the standard model fields. In principle nonrenormalizable couplings could arise from D type Kahler potential terms or F type superpotential terms. However, with conserved R parity, the gauge invariant operators which carry baryon number contain at least 3 standard model fields. A Kahler potential coupling of this type to the decaying particle is dimension 6, but a superpotential coupling to 3 fields is dimension 5. The only invariant made out of 3 standard model fields which carries $U(1)_B$ is $\bar{u}d\bar{d}$. The unique superpotential coupling which satisfies the requirements is therefore

$$W = \frac{\lambda}{M} \phi \bar{u} d\bar{d}$$
(3)

where ϕ is the decaying particle and generation indices are suppressed. Notice that $\bar{u}d\bar{d}$ is odd under *R* parity. So with an unbroken *R* parity (at least) one LSP results from each decay. In addition, if *R* parity is to remain unbroken after the decay, $\phi = 0$ must be the ground state.

Depending on the specific model ϕ might decay through other dimension 5 terms in addition to (3). Decay through superpotential couplings to the other Rodd combinations of 3 standard model fields, namely $L\bar{d}Q$, and $LL\bar{e}$ (which do not carry baryon number) would still give at least one LSP per decay, but dilutes the baryon number (for the decay from a condensate with a particle-antiparticle asymmetry discussed below). In GUT theories, such operators are in general related by GUT symmetries. For example, in SU(5) models $\bar{u}d\bar{d}$, $L\bar{d}Q$ and $LL\bar{e}$ are contained in $\bar{5}\bar{5}10$. The existence of these other decay channels related by SU(5)would dilute the baryon number by a factor $\frac{3}{7}$. All other dimension 5 couplings are through operators which do not carry baryon number. These couplings include: 1) superpotential couplings to R even combinations of standard model fields, namely $QH_u\bar{u}, QH_d\bar{d}$, and $LH_d\bar{e}$, 2) Kahler potential couplings

$$\frac{\lambda'}{M}\phi\chi^{\dagger}\chi\tag{4}$$

4

where χ is a light field, and 3) ϕ dependence of the gauge kinetic functions

$$\frac{g^2}{32\pi^2 M} \phi W^{\alpha} W_{\alpha} \tag{5}$$

where W_{α} is the field strength for a light gauge supermultiplet. All of these decay modes of course do not contribute to the baryon asymmetry. However, if $m_{\phi} < 2m_{\rm LSP}$ no LSPs result either. So if ϕ is light enough these decay modes do not affect the relic $\Omega_b/\Omega_{\rm LSP}$. Finally, the coupling ϕLH_u , if present, would allow decay through a renormalizable operator, giving a very large decay temperature. In addition, it would cause ϕ to pair up with some linear combination of neutrinos after electroweak symmetry breaking, giving a Dirac neutrino with weak scale mass. This (dangerous) coupling must therefore be restricted in some way (as in the toy model given in the next section).

The decay through the operator (3) can in principle lead to a net baryon asymmetry, parameterized by $\epsilon = \langle N_b \rangle / \langle N_{\text{LSP}} \rangle$, where $\langle N_b \rangle$ and $\langle N_{\text{LSP}} \rangle$ are the average number of baryons and LSPs resulting from each decay. In order for $\Omega_b / \Omega_{\text{LSP}} \sim \mathcal{O}(m_b / m_{\text{LSP}})$ to hold, ϵ should not be too small. Direct production of a baryon asymmetry in the decay requires decay channels which carry different baryon number, final state interactions, and CP violating interference terms which contain at least two baryon violating couplings [11]. With conserved R parity, baryon number is violated only by nonrenormalizable operators, giving negligible interference terms. Any baryon asymmetry produced directly in the decay is therefore insignificant [12]. However, a nonzero $\langle N_b \rangle$ will be transferred to baryons through the operator (3) if there is an initial particle-antiparticle asymmetry in the population of decaying particles. Since ϵ should not be too small, there must exist a near maximal asymmetry in the decaying population.

Such a large asymmetry might appear hard to achieve. However, the coherent production of a scalar condensate along a supersymmetric flat direction can give rise to a large asymmetry in the condensate [2,6]. Here, flat direction refers to a direction in field space on which the perturbative potential vanishes at the renormalizable level. Such directions are generic in supersymmetric theories. The nonrenormalization theorem protects these directions from being lifted by quantum corrections [13]. In the presence of SUSY breaking, a potential can arise though. Whether or not a condensate is actually generated along a flat direction depends on the sign of the SUSY breaking soft mass term at early times, $m^2\phi^*\phi$, where ϕ parameterizes the flat direction. When $H \gtrsim m_{3/2}$ the finite energy density of the universe induces soft parameters along flat directions with a scale set by the Hubble constant [6]. If the induced $m^2 > 0$, the origin is stable, and the large expectation values required to form a condensate do not arise. However, if the induced $m^2 < 0$, the origin is unstable and large expectation values can develop. In this case, if the flat direction is lifted at order n by a nonrenormalizable operator in the superpotential

$$W = \frac{\beta}{nM_n^{n-3}}\phi^n \tag{6}$$

then the relevant part of the potential along the flat direction is

$$V(\phi) = (cH^2 + m_{\phi}^2)|\phi|^2 + \left(\frac{(A+aH)\beta\phi^n}{nM_n^{n-3}} + h.c.\right) + |\beta|^2 \frac{|\phi|^{2n-2}}{M_n^{2n-6}}$$
(7)

where $m_{\phi} \sim A \sim m_{3/2}$ are soft parameters arising from hidden sector SUSY breaking, and $c \sim a \sim \mathcal{O}(1)$ are the soft parameters induced by the finite energy density [6]. Here the scale M_n may in general be (much) different than the scale of the operators which allow ϕ to decay. For c < 0 the expectation value along the flat direction is determined at early times by a balance between the mass term and nonrenormalizable terms. If $m_{\phi}^2 > 0$, then when $H \sim m_{3/2}$ the origin becomes stable and the field begins to oscillate freely with a large initial value. However, at just this time since the expectation value of the field is determined by a balance between the mass and nonrenormalizable terms, the U(1) violating A term necessarily has the same magnitude. Depending on the initial phase of the field, the presence of the A term with this magnitude can lead to a near maximal asymmetry in the condensate. So if a condensate is produced along a flat direction which is lifted by a nonrenormalizable superpotential, it naturally has a large asymmetry. For a flat direction made of squark or slepton fields, this is the mechanism of baryogenesis proposed by Affleck and Dine [2]. Here however, the initial condensate asymmetry is in the ϕ field, and is only transferred to baryons by decay through the operator (3).

The final, and perhaps most nontrivial requirement, is that $\Omega_{\text{LSP}} + \Omega_b \simeq 1$. If the decaying particles dominate the energy density at the time of decay, the universe is in a matter dominated era at that epoch. However, since m_{LSP} and m_{ϕ} are the same order this would imply the universe remained matter dominated below this temperature. Matter domination from such an early epoch is incompatible with nucleosynthesis [7]. While it may have been natural for the condensate to dominate the energy density, this is clearly unacceptable [14]. The condensate must have a small enough energy density so that matter domination from the relic LSPs starts at a temperature of $T \sim 5 \Omega h^2$ eV, where $h = H/(100 \text{ km s}^{-1} \text{ Mpc}^{-1})$, and H is the Hubble constant now. Assuming critical density, $N_{\text{LSP}} = 1$ in the decay, and given the current temperature, the condensate number density at the time of decay can be parameterized as

$$\frac{n_{\phi}}{s} \simeq 7 \times 10^{-11} h^2 \left(\frac{50 \text{ GeV}}{m_{\text{LSP}}}\right) \tag{8}$$

where $s = (2\pi^2 g_*/45)T^3$ is the entropy density at the time of decay. This is much less than a thermal number density, $n/s \sim 1/g_*$. Now the total density in the condensate is determined by the expectation value when the field begins to oscillate freely. From (7) the value of the field when oscillations begin $(H \sim m_{3/2})$ is

$$\phi_0 \simeq \left(\frac{\alpha m_\phi M_n^{n-3}}{\beta}\right)^{\frac{1}{n-2}} \tag{9}$$

where α is a constant of order unity. The fractional energy density in the condensate when oscillations begin is $\rho_{\phi}/\rho_{\text{tot}} \simeq \phi_0^2/M_p^2 \ll 1$. In an inflationary scenario with a reheat temperature low enough to avoid overproducing gravitinos by thermal rescatterings, the universe is in an inflaton matter dominated era when $H \sim m_{3/2}$ [6]. So $\rho_{\phi}/\rho_{\text{tot}}$ stays roughly constant until the inflaton decays. After the inflaton decays the condensate density per entropy density is

$$\frac{n_{\phi}}{s} \sim \frac{T_R}{m_{\phi} M_p^2} \left(\frac{m_{\phi} M^{n-3}}{\beta}\right)^{\frac{2}{n-2}} \tag{10}$$

where T_R is the reheat temperature after inflation. Without additional entropy releases n_{ϕ}/s stays constant until the time of decay. So the relic fractional density in the condensate is determined by the order at which the flat direction is lifted, and the reheat temperature after inflation. For $M_n \sim M_p$ and $n \geq 6$ this is generally too large for reasonable reheat temperatures. However for n = 4 and $M_n/\beta \sim M_p$, $n_{\phi}/s \sim T_R/M_p$. With $T_R \sim 10^8$ GeV, this gives just the required condensate density. Therefore if the late decaying condensate arises from coherent production along a direction which is lifted by a Planck suppressed fourth order term in the superpotential, the required relic density naturally arises for reasonable values of the reheat temperature after inflation.

In addition to the required fourth order Planck suppressed term in the superpotential, there are in general higher order SUSY breaking terms (in addition to the mass term) which are suppressed by the Planck scale. Assuming hidden sector SUSY breaking these give a general form for the soft potential of

$$V_s(\phi) = m_{3/2}^2 M_p^2 \mathcal{F}(\phi/M_p)$$
(11)

However, just on energetic grounds the nonrenormalizable term in the superpotential forces $\phi \ll M_p$. The higher order corrections in (11) are therefore unimportant. In addition, there are higher order soft terms generated by integrating out fields which gain mass at the scale M. These are of the general form

$$V_{s}(\phi) = m_{3/2}^{2} M^{2} \mathcal{G}(\phi/M)$$
(12)

With M in the range required to give an acceptable decay temperature for ϕ , these in higher order terms are less important for $H \sim m_{3/2}$ than the terms in (7) with n = 4. The higher order SUSY breaking potential terms for ϕ therefore do not spoil the expectation that the condensate carries a large asymmetry, or the prediction for the relic density.

A TOY MODEL FOR BARYONS AND LSPS FROM LATE DECAY

It is easy to build models which satisfy all the requirements outlined in the previous section. As an existence proof, consider the following toy model. The flat direction required for the coherent production can be parameterized by a singlet field ϕ . In principle this could be a composite field in some sector of the theory, but here will be taken to be an elementary singlet for simplicity. The singlet ϕ should be protected from obtaining a large mass while allowing the operator (3). This can be enforced with discrete symmetries. For example, under a Z_4 discrete R symmetry the superpotential transforms as $W \to -W$ [15]. If ϕ and all the \bar{u} and \bar{d} transform as $f \to e^{i\pi/4}f$, where $f \in (\phi, \bar{u}, \bar{d})$, then the operator (3) is

allowed while a superpotential mass term, $m\phi\phi$, is not allowed. The usual Yukawa couplings, $\lambda_u Q H_u \bar{u}$, $\lambda_d Q H_d \bar{d}$, and $\lambda_e L H_d \bar{e}$, are allowed if the other standard model fields transform under the Z_4 as $L \to L$, $Q \to e^{i\pi/4}Q$, $h \to e^{i\pi/2}h$, where $h \in$ (H_u, H_d, \bar{e}) . The operator (3) must be generated by integrating out particles with intermediate scale mass. This can be accomplished in this model by introducing a Dirac pair U' and \bar{U}' , with mass, $m_U U' \bar{U}'$, and Yukawa couplings $g \bar{U}' \bar{d} \bar{d}$, and $g_\phi \phi \bar{u} U'$. These couplings and Dirac mass can be enforced by the transformation $U' \to e^{i\pi/2}U'$ and $\bar{U}' \to e^{i\pi/2}\bar{U}'$. The mass scale m_U could arise from dynamics which preserves the discrete symmetry. Integrating out the Dirac pair gives the operator (3) with $\lambda/M = gg_{\phi}/m_U$. The operator (5) is also generated at the scale M, but is suppressed by a loop factor compared with (3). So in this model the dominant decay mode is $\phi \to \bar{u} \bar{d} \bar{d}$. Finally, the dangerous superpotential coupling $\phi L H_u$ is restricted by the discrete symmetry.

The flat direction ϕ can be lifted by nonrenormalizable terms in the superpotential. With the Z_4 R symmetry, the lowest order term in the superpotential is ϕ^4 . Such an operator is not generated at the intermediate scale, but presumably can arise directly at the Planck scale

$$W = \frac{\beta}{M_p} \phi^4 \tag{13}$$

And, as discussed in the last section, an operator lifted at fourth order in the superpotential and suppressed by the Planck scale is precisely what is required to give the correct magnitude for the dark matter and baryon densities. In addition to the fourth order term in the superpotential, ϕ is lifted by a fourth order SUSY breaking term in the soft potential generated by integrating out the heavy Dirac pair $U'\bar{U}'$

-

$$V_s(\phi) \simeq \frac{g_\phi^4 m_{3/2}^2}{16\pi^2 m_U^2} (\phi^* \phi)^2 \tag{14}$$

However, as discussed in the previous section, for $H \sim m_{3/2}$, terms of this order are subdominant compared with (13).

Acceptable soft SUSY breaking terms can also result in this model. In order to allow visible sector gaugino masses the $Z_4 R$ symmetry must be broken in the SUSY breaking sector to $Z_2 R$ parity. For definiteness consider a hidden sector scenario

in which SUSY breaking is transmitted by Planck suppressed interactions. The breaking to Z_2 R parity can be accomplished with a hidden sector field z which is invariant under Z_4 and breaks SUSY by an auxiliary component expectation value $\langle F_z \rangle \sim \sqrt{m_{3/2} M_p}$. In addition, the soft A term $A\beta \phi^4/M_p$, required to generate an asymmetry in the ϕ condensate, can arise from supergravity interactions, and the Kahler potential coupling $\frac{1}{M_{\rm p}} \int d^4\theta z \phi^{\dagger} \phi$. Dimension 3 soft A terms for standard model fields arise from similar couplings. A soft H_uH_d scalar mass and weak scale μ term can arise from Kahler potential couplings $\frac{1}{M_{e}^{2}}\int d^{4}\theta z^{\dagger}z'H_{u}H_{d}$ and $\frac{1}{M_p}\int d^4\theta z' H_u H_d$, where z' is a hidden sector field which participates in SUSY breaking and transforms as $z' \rightarrow -z'$ under the discrete symmetry. A weak scale mass for the flat direction, $m_{3/2}^2 \phi^* \phi$, results from supergravity interactions and/or Kahler potential couplings with the hidden sector. However, most importantly, with the hidden sector couplings sketched above, the Z_4 symmetry does not allow a soft mass term $m_{3/2}^2 \phi \phi$ from Kahler potential couplings, which would violate the U(1) carried by ϕ . The classical evolution of the condensate at late times therefore preserves the asymmetry generated when the coherent oscillations begin.

So in this model all the requirements are satisfied with a single discrete symmetry. Although the model is perhaps unrealistically simple, it demonstrates that all the requirements for baryons and LSPs from late decay of a condensate can be met in a technically natural manner.

CONCLUSIONS

In the late decay scenario outlined here, the baryon and LSP densities are related by

$$\frac{\Omega_b}{\Omega_{\rm LSP}} \simeq \epsilon \; \frac{m_b}{m_{\rm LSP}} \tag{15}$$

As discussed above, $\epsilon < 1$ if one LSP results from each decay. Under the assumption that the total density is near critical, $\Omega_b + \Omega_{\text{LSP}} \simeq 1$, a lower limit on the baryon density then gives an upper limit on the LSP mass in this scheme. For $\Omega_b \ll \Omega_{\text{LSP}}$,

$$m_{\rm LSP} \simeq \epsilon \; \frac{m_b}{\Omega_b}$$

The absolute lower bound on Ω_b comes from the observed density of luminous matter, $\Omega_b \gtrsim .007$. This gives an upper limit of $m_{\text{LSP}} \lesssim 140 \epsilon$ GeV. A more

stringent upper limit comes from nucleosynthesis. The primordial light element abundances depend on the baryon to entropy ratio at the time of nucleosynthesis. Comparison of the calculated and observed abundances gives upper and lower limits on the baryon density, $.01 \leq \Omega_b h^2 \leq .015$ [16]. The nucleosynthesis lower limit on the baryon density gives the upper limit $m_{\rm LSP} \leq 100h^2 \epsilon$ GeV. So the LSP is expected to be fairly light in this late decay scenario. For example, in SU(5)models for which $\epsilon < \frac{3}{7}$, with a hubble constant h < .8, the upper limit on the LSP mass is $m_{\rm LSP} < 30$ GeV.

In addition to the LSPs arising from the late decaying condensate there will be a population of LSPs arising from thermal freeze out. However, the low mass required for the late decay scenario can give a freeze out density which is well below critical. For an LSP which has a sizeable mixture of Higgsino and gaugino components, annihilation through s-channel Z exchange is very efficient and leaves a very small relic density from freeze out [17,18]. For a mostly Higgsino LSP, coannihilation with the other Higgsino states also leads to negligible relic density [19]. A small relic density for a light nearly pure gaugino LSP can also result from annihilation through t-channel squark and slepton exchange if one of the sleptons or squarks are light [18]. So depending on the precise composition of the LSP, the late decay can give the dominant contribution to the relic LSP density. Independent of the production of a baryon asymmetry, late decay is an interesting source of relic LSPs in the low mass regime. In fact, if the LSP was found to be in a region of parameter space for which the freeze out density was too small to give closure (such as the light Higgsino or mixed Higgsino-gaugino regions) the only alternate source for LSPs would be a late decay below the freeze out temperature.

In conclusion, the ratio of dark matter to baryon density can naturally be $\mathcal{O}(10-100)$ if stable weak scale mass particles and baryons result in roughly equal amounts from the late decay of a particle. A natural way in which this can occur is for a condensate with a large net asymmetry to decay to R odd combinations of standard model fields. Supersymmetric theories with a conserved R parity can in principle have all the ingredients to realize this scenario.

I would like to thank R. Brandenberger, M. Dine, J. Primack, L. Randall, and U. Sarid for useful discussions, and J. Frieman, R. Malaney, and G. Steigman for discussions about nucleosynthesis.

REFERENCES

- 1. For a review of electroweak baryogenesis see A. Cohen, D. Kaplan, and A. Nelson, Ann. Rev. Nucl. Part. Sci. 43 (1993) 27.
- 2. I. Affleck and M. Dine, Nucl. Phys. B 249 (1985) 361. For a recent discussion of the Affleck-Dine mechanism of baryogenesis see ref. [6].
- For a review of baryogenesis from the out of equilibrium decay of heavy particles see E. Kolb and M. Turner, Ann. Rev. Nucl. Part. Sci. 33 (1983) 645.
- S. Nussinov, Phys. Lett. B 165 (1985) 55; R. Chivukula and T. Walker, Nucl. Phys. B 329 (1990) 445; S. Barr, R. Chivukula, and E. Farhi, Phys. Lett. B 241 (1990) 387; D. Kaplan, Phys. Rev. Lett. 68 (1992) 741.
- 5. D. Lyth and E. Stewart, preprint Lancaster-TH-9502, hep-ph/9502417.
- M. Dine, L. Randall, and S. Thomas, preprint SLAC-PUB-95-6776, hepph/9503303, to appear in Phys. Rev. Lett.; M. Dine, L. Randall, and S. Thomas, preprint SLAC-PUB-95-6846.
- D. Lindley, Ap. J. 294 (1985) 1; J. Ellis, D. Nanopoulos, and S. Sarkar, Nucl. Phys. B 259 (1985) 175; S. Dimopolous, R. Esmailzadeh, L. Hall, and G. Starkman, Nucl. Phys. B 311 (1988) 699; J. Ellis, G. Gelmini, J. Lopez,
- D. Nanopoulos, and S. Sarkar, Nucl. Phys. B 373 (1992) 399.

-

- M. Reno and D. Seckel, Phys. Rev. D 37 (1988) 3441; G. Lazarides, R. Schaefer, D. Seckel, and Q. Shafi, Nucl. Phys. B 346 (1990) 193.
- J. Ellis, A. Linde, and D. Nanopoulos, Phys. Lett. B 118 (1982) 59; D. Nanopoulos, K. Olive, and M. Srednicki, Phys. Lett. B 127 (1983) 30; M.Yu. Khlopov and A.D. Linde, Phys. Lett. 138B (1984) 265; J. Ellis, J.E. Kim and D.V. Nanopoulos, Phys. Lett. 145B (1984) 181.
- C. Coughlan, W. Fischler, E. Kolb, S. Raby, and G. Ross, Phys. Lett. B 131 (1983) 59; J. Ellis, D.V. Nanopoulos and M. Quiros, Phys. Lett. B 174 (1986) 176.
- 11. D. Nanopoulos and S. Weinberg, Phys. Rev. D 20 (1979) 2484.
- 12. A baryon symmetric but CP asymmetric primary decay can be transformed into a non-negligible baryon asymmetry in the decay chain by renormalizable baryon violating interactions. However, with minimal low energy particle content this requires R parity violation. S. Dimopoulos and L. Hall, Phys.

Lett. B 196 (1987) 135; J. Cline and S. Raby, Phys. Rev. D 43 (1991) 1781;
S. Mollerach and E. Roulet, Phys. Lett. B 281 (1992) 303.

- M. Grisaru, W. Siegal, and M. Rocek, Nucl. Phys. B 159 (1979) 429; N. Seiberg, Phys. Lett. B 318 (1993) 469.
- The relic LSP density arising from gravitino decay has been considered by L. Krauss, Nucl. Phys. B 227 (1983) 556; and from axino decay by K. Rajagopal, M. Turner, and F. Wilczek, Nucl. Phys. B 358 (1991) 447.
- 15. Under a discrete $Z_N R$ transformation, a chiral superfield transforms as $\Phi \to e^{2\pi i R/N} \Phi$, where R is the $Z_N R$ charge of the field. The scalar, fermionic, and auxiliary components transform as $A \to e^{2\pi i R/N} A$, $\psi \to e^{2\pi i (R-1)/N} \psi$, and $F \to e^{2\pi i (R-2)/N} F$. The superpotential transforms as $W \to e^{4\pi i/N} W$.
- 16. T. Walker, G. Steigman, D. Schramm, K. Olive, and H. Kang, Astrophys. J. **376** (1991) 51; W. Smith, L. Kawano, and R. Malaney, Astrophys. J. Suppl. **85** (1993) 219; For recent discussions of the limits on $\Omega_b h^2$ from nucleosynthesis see P. Kernan and L. Krauss, Phys. Rev. Lett. **72** (1994) 3309; K. Olive and G. Steigman, preprint UMN-TH-1230/94, OSU-TA-6/94, astro-ph/9405022.
- 17. J. Ellis, D. Nanopoulos, L. Roszkowski, and D. Schramm, Phys. Lett. B 245 (1990) 251.
- L. Roszkowski, Phys. Lett. B 262 (1991) 59; L. Roszkowski, Phys. Lett. B
 278 (1992) 147; J. McDonald, K. Olive, and M. Srednicki, Phys. Lett. B
 283 (1992) 80; A. Bottino, V. De Alfaro, N. Fornengo, G. Mignola, and M. Pignone, Astropart. Phys. 2 (1994) 67.
- K. Griest and D. Seckel, Phys. Rev. D 43 (1991) 3191; S. Mizuta and M. Yamaguchi, Phys. Lett. B 298 (1993) 120; M. Dress and M. Nojiri, Phys. Rev. D 47 (1993) 376.

13