Hypercharge and the Cosmological Baryon Asymmetry.

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Abstract

Stringent bounds on baryon and lepton number violating interactions have been derived from the requirement that such interactions, together with electroweak instantons, do not destroy a cosmological baryon asymmetry produced at an extremely high temperature in the big bang. While these bounds apply in specific models, we find that they are generically evaded. In particular, the only requirement for a theory to avoid these bounds is that it contain charged particles which, during a certain cosmological epoch, carry a non-zero hypercharge asymmetry. Hypercharge neutrality of the universe then dictates that the remaining particles must carry a compensating hypercharge density, which is necessarily shared amongst them so as to give a baryon asymmetry. Hence the generation of a hypercharge density in a sector of the theory forces the universe to have a baryon asymmetry.

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1 Introduction

Various authors [1, 2] have placed cosmological bounds on the size of baryon and lepton number violating interactions in theories where baryogenesis occurs before the electroweak phase transition. The baryon asymmetry of the universe is threatened by a combination of these interactions and a large electroweak instanton rate [3, 4]. Electroweak instanton interactions are expected to be in equilibrium for temperatures above $T_{\text{min}}$, approximately the weak breaking scale, up to some very high temperature $T_{\text{max}} \simeq 10^{12}$GeV. Such reactions create $SU(2)_L$ transforming fermions out of the vacuum [3]. Lepton and baryon violating interactions such as R-parity breaking terms in supersymmetry [5] or Majorana neutrino masses, when in equilibrium simultaneously with instanton reactions can break all linear combinations of conserved quantum numbers which involve baryon number. Naively, one is led to believe that the baryon asymmetry of the universe is therefore washed away. In this paper we examine the general circumstances in which this outcome is avoided. We find that in many models there will be additional symmetries and, even though these symmetries apparently have nothing to do with baryon number, they automatically lead to a protection of it.

It is well known that a symmetry which involves baryon number itself, such as $B-3L_i$, can preserve the baryon asymmetry [4]. Approximate symmetries involving $B$ have been found in the minimal supersymmetric standard model which can be used to help prevent erasure of the baryon asymmetry [3]. We have found that the protection of the baryon asymmetry is extremely common and is a typical feature of theories with extra symmetries, even when those symmetries do not transform quarks. We illustrate this by a very simple example: assume that there exists a particle, $X$, which carries hypercharge but not $SU(2)$ or $SU(3)$ gauge interactions. Assume that reactions occurring at temperatures well above $T_{\text{min}}$ generate an asymmetry in the $X$ species, and that at lower temperatures the reactions which change $X$ number are sufficiently weak that this $X$ asymmetry persists. A crucial role is played by the requirement that the early universe is hypercharge neutral. Because $X$ particles carry hypercharge, the asymmetry in their number contributes to the hypercharge density of the universe. The remaining particles in the theory must carry an opposite hypercharge density to cancel this. Chemical equilibrium equations specify how this hypercharge density is shared. A baryon asymmetry can develop either through added $B$ violating interactions or once the weak instanton becomes effective. In general any $X$ asymmetry together with chemical equilibrium requires a baryon asymmetry[4]. This illustrates just how easy it is to preserve the baryon asymmetry and, to our way of thinking, puts the issue of direct detection of baryon and lepton number violation back where it belongs: with the experimentalists.

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3 Cline, Kainulainen, and Olive have recently employed the right handed electron as the $X$ particle [7]. However, in this case non-zero Yukawa interactions limit the temperature range over which the right handed electron can maintain a hypercharge asymmetry.

4 Implicit in this discussion is the assumption that the universe is homogeneous.
2 General condition for survival of a baryon asymmetry.

In this section we discuss, in a very general way, the conditions under which an extra $U(1)$ symmetry preserves the cosmological baryon asymmetry.

In thermodynamic equilibrium the number density of particle species $i$ is determined by its chemical potential, $\mu_i$. If a given reaction, say $p_1 + p_2 \rightleftharpoons p_3 + p_4$, is in equilibrium then $\mu_1 + \mu_2 = \mu_3 + \mu_4$. It is straightforward, yet tedious, to solve all chemical equilibrium equations. One can simplify the process by noticing that these equations are the same equations one would write down to determine the $U(1)$ symmetries of the equilibrium theory. One need only replace $\mu_i$ with $q_i$, the charge of particle $i$. Solving for $q_i$ determines the possible assignments of $U(1)$ charge to each particle so that all equilibrium reactions conserve that charge. In general such $U(1)$ symmetries need not be exact symmetries of the Lagrangian. They are symmetries of those interactions in thermal equilibrium at temperature $T$, and we refer to them as effective $U(1)$ symmetries at this temperature.

Thus, a solution to the chemical equilibrium equations corresponds to an assignment of effective $U(1)$ charges to each particle, and the possible effective $U(1)$s in a given theory are usually easy to identify. Suppose that at a certain temperature there are $N$ such effective $U(1)$s: $U(1)_A, A = 1, \ldots, N$, then the most general solution is

$$\mu_i = \sum_A C_A^i q_i^A$$

(1)

where $q_i^A$ is the charge of particle $i$ under $U(1)_A$. The constant $C_A^i$ we refer to as the asymmetry constant for $U(1)_A$. As soon as some interaction which violates $U(1)_A$ comes into thermal equilibrium, $C_A^i$ rapidly tends to zero: $U(1)_A$ is no longer able to support particle asymmetries.

This general solution is restricted, however. We assume that the universe is homogeneous and that no charge asymmetry has developed for the unbroken gauged $U(1)$s of the theory. This forces the charge density for these $U(1)$s to zero. We can write the charge density for $U(1)_A$ as follows,

$$Q_A = \sum_i q_i^A n_i$$

(2)

where $n_i$ is the particle asymmetry density of species $i$. If particle asymmetry densities are small then they can be written, for $T \gg m_i$, as

$$n_i \simeq \frac{T^2}{6 \tilde{g}_i \mu_i}$$

(3)

where $\tilde{g}_i$ is the number of internal degrees of freedom of particle $i$, $g_i$, multiplied by a factor of two for bosons. (However, see reference [8] for an interesting look at small mass.
effects.) Under these conditions the charge density constraint is a simple linear equation in the $\mu_i$s. $Q^A$ can be written using $n_i$ from (3) and $\mu_i$ from (1):

$$Q^A \simeq \frac{T^2}{6} \sum_B C^B \sum_i \tilde{g}_i q_i^B q_i^A = \frac{T^2}{6} \sum_B C^B \overline{B} \cdot \overline{A}$$

(4)

where we define $\overline{B} \cdot \overline{A}$ by

$$\overline{B} \cdot \overline{A} = \sum_i \tilde{g}_i q_i^B q_i^A.$$

(5)

Should the diagonal generators of non-Abelian gauge groups, such as $T_3L$, be included in the list of effective $U(1)$s? The answer is no, as can be seen easily from the above equations. Call such a generator $\alpha$, then neutrality of the universe with respect to this charge requires

$$\sum_i C^A \overline{A} \cdot \overline{\alpha} = 0$$

(6)

where

$$\overline{A} \cdot \overline{\alpha} = \sum_i \tilde{g}_i q_i^A q_i^\alpha.$$

(7)

When $A$ refers to a $U(1)$ generator (not embedded in a non-Abelian gauge group) then $\overline{A} \cdot \overline{\alpha} = 0$. This is because the $\tilde{g}_i$ and $q_i^\alpha$ are the same for all components of an irreducible representation of $\alpha$, and hence the sum in (7) can be written as a sum of zero terms, one for each irreducible multiplet of $\alpha$. When $A = \beta$ is a diagonal generator of a non-Abelian group the orthogonality property of the generators within each multiplet ensures that $\sum_i q_i^\alpha q_i^\beta$ vanishes for $\beta \neq \alpha$. Hence the sum in (6) just has one term: $C^\alpha \overline{\alpha} \cdot \overline{\alpha} = 0$. Since $\overline{\alpha} \cdot \overline{\alpha} \neq 0$, we have proved that $C^\alpha = 0$ follows from $Q^\alpha = 0$. This implies that such $U(1)$s need not be included in the list of effective $U(1)$s.

Now let’s apply this formalism. We are interested in the situation in which additional particles and interactions have been added to the standard model such that at temperatures $T$, $T_c < T < T_{max}$, where $T_c$ is the weak breaking temperature, there are just two effective $U(1)$s: $Y$ and $X$, where $Y = 2(Q - T_3)$ denotes hypercharge and $X$ is an ungauged effective symmetry. The charge neutrality condition (6) when applied to hypercharge gives

$$C^Y = -\frac{X}{Y} \cdot Y.$$  

(8)

Using (8) in equation (4) the asymmetry in baryon number is just

$$n_B \simeq \frac{T^2}{6} C^X \left( \frac{X}{Y} - \frac{X \cdot Y}{Y^2} \right) \cdot \overline{B}.$$  

(9)

where we have rewritten $Q^B$, the baryon density, as $n_B$. This is the general result of this paper. Any effective $U(1)_X$, whether it contains a piece of baryon number or not, will in general contribute to $n_B$ if $C^X \neq 0$. The extension of (9) to many extra $X$ symmetries
is straightforward. Providing such a $U(1)_X$ exists, there is no limit to how large the $B$ and $L$ violating interactions can be.

We will examine the case in which $X$ particles carry no baryon number themselves. Then

$$n_B \simeq \frac{T^2}{6} C^X \left( -\frac{\mathbf{Y} \cdot \mathbf{B}}{Y^2} \right) \mathbf{X} \cdot \mathbf{Y}. \quad (10)$$

In the standard model $\mathbf{Y} \cdot \mathbf{B}/Y^2 = \frac{1}{11}$. Additional particles will change this, but would generally give some non-zero value which we call $\alpha$. Then $n_B \simeq -\frac{T^2}{6} \alpha C^X (\mathbf{X} \cdot \mathbf{Y})$. Thus to obtain $n_B \neq 0$ we require that some particles with $X_i \neq 0$ have $Y_i \neq 0$. Hypercharge neutrality then forces other particles to have an asymmetry, some of which carry baryon number, thus providing a baryon asymmetry.

Cline et al. [7] point out that in the standard model right handed electron number is conserved down to a temperature of about $10^{-4} TeV$, and thus can insure that baryon/lepton violating interactions don’t wash away the baryon asymmetry above this temperature. Thus the standard model already contains $X$ particles in the form of right handed electrons. In section 3, we discuss another possibility, an $X$ symmetry which does not transform any standard model particles. In this case (8) can be rewritten in terms of the hypercharge density carried by the standard model sector, $Q^Y(SM)$, and by the $X$ sector, $Q^Y(X) \equiv \sum_i q_i^Y n_{X_i}$.

$$Q^Y(SM) = Q^Y(X) \quad (11)$$

In terms of $Q^Y(X)$ equation (10) becomes

$$n_B \simeq -\frac{1}{11} Q^Y(X). \quad (12)$$

(We have assumed that $T < 10 TeV$ so that right handed electrons are in equilibrium.)

Equation (12) doesn’t assume that $X$ number density is small or proportional to its chemical potential. Thus it is valid even when the temperature drops below the mass of certain $X$ particles. When this happens the heavier species carrying $X$ might decay into lighter ones. Nevertheless, providing the particles with $X \neq 0$ possess a hypercharge asymmetry the baryon asymmetry will survive. In particular the $X \neq 0$ particles must continue to carry such an asymmetry until a temperature $T_0$, beneath which $B$ and $L$ violating reactions are sufficiently weak that a symmetry having a baryon number component has become an effective $U(1)$. The resulting baryon asymmetry after $X$ decay depends on the specifics of the model. In the least complicated scenario, in which baryon number is a good symmetry below $T_0$, today’s baryon asymmetry is simply derived form (12) and entropy considerations.

We note that it is not necessary for our $X$ sector to be neutral under $SU(2)$. Adding additional $SU(2)$ transforming fermions to the standard model will mean that these particles also take part in instanton mediated reactions. Nevertheless, in a consistent theory, instanton reactions will conserve the hypercharge asymmetry carried by the $X$
sector of the theory. This is true because instantons neither violate hypercharge in the standard model sector nor in the theory overall, and thus must conserve hypercharge in the $X$-sector.

In this section we have tacitly assumed that some component of baryon number is a good symmetry below $T_{c}$, the weak breaking temperature. If this is not the case, then, for temperatures $T$, $T_{0} < T < T_{c}$, the role of hypercharge is played by electric charge. In this case the $X$ sector must carry an electric charge asymmetry.

An intriguing possibility exists if the lightest $X$ particle is stable and electrically neutral. If this is the case, the particle is a candidate for the dark matter in the universe \cite{9,10}. To realize such a scenario, the $X$ sector would still have to maintain a hypercharge asymmetry for temperatures above $T_{0}$. (For convenience, we have assumed $T_{0} \geq T_{c}$.) However, at a lower temperature, charged $X$ particles would decay to standard model particles plus these electrically neutral $X$ particles. If $\Omega_{X}$ is the fraction of the critical density contributed by the electrically neutral $X$ particles then their mass is given by

$$\frac{m_{X}}{m_{\text{proton}}} \approx \frac{\langle q_{X} \rangle}{11} \times 10^{2} \Omega_{X}$$

where $\langle q_{X} \rangle$ is the appropriate average of $X$-particle hypercharges. Low-background Ge detector experiments \cite{11,12} indicate that an electrically neutral dark matter particle with nonzero hypercharge must have a mass greater than $\sim 1000$ GeV. Thus, we can effectively rule out a dark matter $X$ particle with nonzero hypercharge. One possible candidate is the neutral component of a new hyperchargeless $SU(2)$ multiplet. Such a particle is expected to interact via loop diagrams with nuclei and thus its cross section with $Ge$ is approximately $10^{-35} cm^{2}$ or smaller \cite{10}, effectively evading relevant experimental limits \cite{11}. Another candidate is a new particle with no gauge interactions whatsoever \cite{9}.

## 3 A Simple Model.

In this section we illustrate the general ideas discussed above with a very simple model. We add to the standard model a single fermion $X$, of mass $m_{X}$, which is $SU(2)$ neutral but has three units of electric charge. It is unstable, decaying to three charged leptons via the effective interaction

$$\frac{1}{2} \frac{1}{M^{2}} \sum_{ijk} f_{ijk} \left( X e_{R}^{i} \right) \left( (e_{R}^{j})^{T} C e_{R}^{k} \right) + \text{h.c.}$$

where $e_{R}^{i}$ is the right handed lepton field of flavor $i$, $X$ is the $X$ particle field, $C$ is the charge conjugation matrix, $M$ is a constant with units of energy, and $f_{ijk} (= f_{ikj})$ is a flavor dependent constant of order 1. In addition we let our model include unspecified lepton and/or baryon violating terms which together with the instanton reaction break all linear combinations of $B$ and $L$ numbers.
Both the mass of the $X$ particle, $m_X$, and the constant $M$ are constrained by the various requirements of our theory. First we must insure that the $X$ asymmetry develops before all baryon violating interactions fall out of equilibrium. Otherwise the $X$ asymmetry has no effect on baryon number. Let $T_X$ be the temperature at which $X$ violating reactions drop out of equilibrium. Without specifying the exact scenario, we assume that an $X$ asymmetry develops at some temperature lower than $T_X$ but above the temperature at which instantons freeze out (See [13] and references therein for numerous methods by which number asymmetries can develop). In this way the instanton reaction provides the baryon violation required for our mechanism to work. This is a convenient choice, but not a necessary one if other baryon violation exists in the theory.

It is interesting to note that the only baryon violation required in this model is instantons. If an $X$ asymmetry exists or develops during the epoch in which instantons are in equilibrium then it will necessarily generate a proportional baryon asymmetry.

In our example $X$ particles will eventually decay into standard model particles. Various constraints must be imposed on this decay. To make things simple we require $X$ particles to survive past the temperature at which instantons freeze out. We assume that after this temperature baryon number is a good symmetry. Thus, the only possible effect on the produced baryon density comes from the change in entropy of the universe upon $X$ decay.

The standard nucleosynthesis scenario places limits on this decay [13]. If $X$ particles decay after nucleosynthesis, they must not dump more than a factor of $\sim 15$ times the entropy density present at the time of nucleosynthesis. If they did than the observed baryon to photon density would be incompatible with standard nucleosynthesis. Also, if the mass of the $X$ particle is larger than a few $MeV$, which it must be to avoid strict limits on the width of the $Z$ boson, then energetic photons from $X$ decay can destroy too much deuterium. Further, depending on the era of decay, photons from $X$ decay can destroy the uniformity of the cosmic microwave background radiation or contribute too much to the diffuse photon background. If $X$ particles decay before nucleosynthesis, their mass and density prior to decay must be compatible with the known baryon to photon ratio, $\eta$, during nucleosynthesis.

Let us examine our first constraint. The rate for $X$ violating 4-fermion interactions is given by

$$\Gamma_X \simeq \frac{49 f^2 \pi^5}{12960 \zeta(3)} \frac{T^5}{M^4}$$

where $f^2$ is an average of terms like $f_{ijk} f_{lmn}$, and we have dropped terms of order $\frac{m_X}{T}$.\footnote{For convenience, and because we are interested in the order of magnitude of our results, we assume $g_* \simeq 106$ independent of the temperature.}

The Hubble constant, $H$, is $17 \frac{T^2}{M_p}$. The 4-fermion interaction drops out of equilibrium when its rate falls below the Hubble expansion rate\footnote{For convenience, and because we are interested in the order of magnitude of our results, we assume $g_* \simeq 106$ independent of the temperature.}. Calling the temperature at which this occurs $T_X$, we have
\[ M^4 \simeq \frac{49f^2\pi^5}{220320\zeta(3)} M_\pi T_X^3. \] (13)

Although \( X \) number changing interactions freeze out at \( T_X \), \( X \) particles stay in thermodynamic equilibrium below this temperature through their gauge interactions. These gauge interactions freeze out at a much lower temperature given by the standard cold relic freeze out criteria.

Now we examine the decay of the \( X \) particles. The decay rate for these particles is given by

\[ \Gamma \simeq \frac{f^2}{256\pi^3} \frac{m_X^5}{M^4} \]

where we have ignored terms of order the temperature over \( m_X \) since they will be seen to be negligible. The \( X \) particles decay when this rate is approximately equal to the Hubble expansion rate. Calling the temperature at which these rates become equal \( T_D \), we have

\[ M^4 \simeq \frac{f^2}{4352\pi^3} \frac{m_X^5}{T_D^2} M_\pi. \] (14)

If significant entropy is generated by \( X \) decay then \( T_D \) is the “reheat” temperature after decay.

Equations (13) and (14) can be combined to give

\[ m_X^5 \simeq 7.6 \times 10^3 T_X^3 T_D^2 \]

In figure 1 we plot the allowed parameter space by considering the constraints discussed above. (We have assumed \( T_D \leq T_{\text{min}} \simeq 10^2 \text{GeV} \) and required \( T_X > 10T_{\text{min}} \).

The diagonal dotted lines in this figure are lines of constant \( T_X \) and are labeled in GeV. The allowed region is divided up into three regimes. The first, corresponding to \( T_D > 10^{-3} \text{GeV} \), covers the case in which \( X \) particles decay before the onset of nucleosynthesis. In this case \( X \) density just before decay may be quite large, leading to an early matter dominated era and a significant increase in entropy density upon \( X \) decay. This is because for large \( m_X \), \( X \) particle gauge interactions freeze out when there is still a large anti-\( X \) particle density. In this situation, the \( X \) number asymmetry is a small fraction of the symmetric relic freeze out density. A large symmetric relic density leads to large entropy dumping when \( X \) particles decay. Let us call the factor by which entropy is increased \( R \). Since, in our model, today’s observed baryon asymmetry is proportional to the \( X \) asymmetry divided by \( R \), a large \( X \) asymmetry is required when \( R \) is large. We have plotted a dot-dashed line which corresponds to the onset of significant entropy generation when \( X \) particles decay. At this line entropy is increased by 10\% upon \( X \) decay. As we rise above this line the amount of entropy generated when the \( X \) particles decay increases. At the top boundary of our allowed region the \( X \) asymmetry required to generate today’s observed baryon asymmetry becomes infinite. Above this line there is no way to generate enough baryon asymmetry.

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In the second regime $10^{-4} \text{GeV} < T_D < 10^{-3} \text{GeV}$, during which nucleosynthesis is taking place we impose the conservative requirement that $X$ decay increases the universe’s entropy by less than 10%. This is shown as a dip in the top boundary of the allowed region.

The last regime, $T_D < 10^{-4} \text{GeV}$, in which $X$ particles decay after nucleosynthesis, is bounded on the left by the requirement that decay products don’t destroy too much deuterium [14]. The curved line marked with an arrow takes account of this limit (We have used Lindley’s rough calculation for heavy dark matter particles [14]). This constraint is more severe than those arising from cosmic microwave background and diffuse photon background observations. The top limit of this region is determined by entropy dumping considerations. Since, in this case $X$ particles are still present during nucleosynthesis, we know that the required $X$ asymmetry is equal to $-\frac{3x}{11}$ times the baryon asymmetry at the time of nucleosynthesis. When $X$ particles decay they can increase the entropy and thus decrease the value of $\eta$ today relative to its value during nucleosynthesis. We allow at most a decrease by a factor of 15, and this gives us our top limit. Figure 1 illustrates how general our mechanism is. The $X$ particle’s mass can range over 12 orders of magnitude, from $45 \text{GeV}$ to $10^{12} \text{GeV}$.

4 Conclusion

We have shown that in order to avoid the strict cosmological limits placed on lepton and baryon number violating interactions it is not necessary to resort to low temperature baryon generation or to the addition of new symmetries which affect baryons. Any symmetry which allows one sector of the theory to acquire a net hypercharge density will suffice. This includes a symmetry under which standard model particles are neutral, as our example shows. The key observation is that, although this new symmetry seems decoupled from the rest of the theory, the gauged $U(1)$ symmetries can connect it. Thus an asymmetry in $X$ particles, because they are charged, is enough to ensure a proportional asymmetry in all charged particles independent of whether their particle number is conserved or not.

If a scenario similar to the one proposed here was realized in the early universe, than experimental searches for lepton and baryon violating interactions may prove successful. Such a success would not only directly signal exciting new $L$ and/or $B$ number violating physics, but would also indirectly signal the existence of a baryon number protection mechanism.

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FIGURE CAPTION

Figure 1: The allowed parameter space in our example is shown, bounded by solid lines. We have assumed $T_D \leq T_{\text{min}} \simeq 10^2 \text{GeV}$ and required $T_X > 10 T_{\text{min}}$ and $m_X > 45 \text{ GeV}$. The diagonal dotted lines are lines of constant $T_D$, and are labeled in $\text{GeV}$. Our parameter $M$ is also constant on these dotted lines, $M = 2.9 \times 10^4 \left( \frac{T_X}{\text{GeV}} \right)^{3/4} \text{ GeV}$. On the dot-dashed line the entropy of the universe is increased by 10% when $X$ particles decay. In determining this line as well as the top boundary line we have assumed that $X$ particle gauge interactions freeze out according to the standard cold relic freeze out criteria [13]. We have made conservative assumptions in determining the relative increase in entropy upon $X$ decay; allowing the cosmic scale factor to scale as $t^n$ where $n$ ranges from $1/2$ to $2/3$. We have used a value for $\eta$ at the time of nucleosynthesis equal to $\left( \frac{11}{4} \right) 3 \times 10^{-10}$.
References

[1] M. Fukugita and T. Yanagida, Phys. Rev. D 42, 1285 (1990); B. Campbell, S. Davidson, J. Ellis and K.A. Olive, Phys. Lett. B 256, 457 (1991); Astropart. Phys. 1, 77 (1992); J. A. Harvey and M. S. Turner, Phys. Rev. D 42, 3344 (1990).

[2] A.E. Nelson and S.M. Barr, Phys. Lett. B 246, 141 (1991); W. Fischler, G.F. Giudice, R.G. Leigh and S. Paban, Phys. Lett. B 258, 45 (1991).

[3] G. ’t Hooft, Phys. Rev. Lett. 37, 8 (1976); Phys. Rev. D 14, 3432 (1976).

[4] V. Kuzmin, V. Rubakov, M. Shaposhnikov, Phys. Lett. B 155, 36 (1985); P. Arnold, L. McLerran, Phys. Rev. D 37, 1020 (1988); N. S. Manton, Phys. Rev. D 28, 2019 (1983); F. R. Klinkhammer and N. S. Manton, Phys. Rev. D 30, 2212 (1984).

[5] L.J. Hall and Suzuki, Nucl. Phys. B 231, 419 (1984); C. Aulakh and R. Mohapatra, Phys. Lett. B 119, 136 (1983); I.H. Lee, Nucl. Phys. B 246, 120 (1984); G.G. Ross and J.W.F. Valle, Phys. Lett. B 151, 375 (1985); S. Dawson, Nucl. Phys. B 261, 297 (1985); F. Zwirner, Phys. Lett. B 132, 103 (1983); R. Barbieri, and A. Masiero, Nucl. Phys. B 267, 679 (1986); S. Dimopoulos and L. J. Hall, Phys. Lett. B 196, 135 (1987).

[6] L. E. Ibáñez and Fernando Quevedo, Phys. Lett. B 283, 261 (1992).

[7] J. Cline, K Kainulainen, K. A. Olive, University of Minnesota preprint, UMN-TH-1201/93 (1993);

[8] H. Dreiner and G.G. Ross, University of Oxford preprint OUTP-92-08P.

[9] S.M. Barr, Phys. Rev. D 44, 3062 (1991); D.B. Kaplan, Phys. Rev. Lett. 68, 741 (1992).

[10] S. Dodelson, B.R. Greene, L.M. Widrow, Nuc. Phys. B 372, 467 (1992).

[11] D.O. Caldwell et al., Phys. Rev. Lett. 61, 510 (1988).

[12] S.P. Ahlen et al., Phys. Lett. B 195, 603 (1987).

[13] E.W. Kolb and M. S. Turner, The Early Universe (Addison-Wesley, Redwood City, CA, 1990) and references therein.

[14] D. Lindley, Astrophys. J. 294, 1 (1985); J. Audouze, D. Lindley and J. Silk, Astrophys. J. 293, L53 (1985).

[15] J. Ellis, G. Gelmini, C. Jarlskog, G.G. Ross and J.W.F. Valle, Phys. Lett. B 150, 142 (1985).
[16] S. M. Barr, R. S. Chivukula and E. Fahri, Phys. Lett. B 241, 387 (1991); S. M. Barr, Phys. Rev. D 44, 3062 (1992); D. Kaplan, Phys. Rev. Lett. 68, 741 (1992).
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