Post–Sphaleron Baryogenesis

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Abstract

We present a new mechanism for generating the baryon asymmetry of the universe directly in the decay of a singlet scalar field $S_r$ with a weak scale mass and a high dimensional baryon number violating coupling. Unlike most currently popular models, this mechanism, which becomes effective after the electroweak phase transition, does not rely on the sphalerons for inducing a nonzero baryon number. CP asymmetry in $S_r$ decay arises through loop diagrams involving the exchange of $W^\pm$ gauge bosons, and is suppressed by light quark masses, leading naturally to a value of $\eta_B \sim 10^{-10}$. We show that the simplest realization of this mechanism, which uses a six quark $\Delta B = 2$ operator, predicts colored scalars accessible to the LHC, and neutron–antineutron oscillation within reach of the next generation experiments.
1 Introduction

Recent developments in particle physics have had profound impact on cosmology. One of the most far-reaching consequences has been the possibility that new interactions beyond the standard model can explain the origin of matter–antimatter asymmetry of the universe as a dynamical phenomenon. There are currently several attractive scenarios which achieve this, the two most widely discussed ones being (i) baryogenesis via leptogenesis [1], which is connected to the seesaw mechanism and neutrino masses, and (ii) weak scale baryogenesis [2], which involves supersymmetric or multi–Higgs extensions of the standard model. Both these proposals depend crucially on the properties of the electroweak sphaleron [3] which serves as the source of $B$ violation. Since the nature of new physics beyond the standard model remains unknown presently, it is important to explore alternative mechanisms that can explain the matter–antimatter asymmetry while yielding testable consequences. In this letter we suggest and explore one such alternative.

The salient feature of our proposal is that baryogenesis occurs via the direct decay of a scalar boson $S_r$ having a weak scale mass and a high dimensional baryon violating coupling. $S_r$ is the real part of a baryon number carrying complex scalar $S$, which acquires a vacuum expectation value (vev). The decays $S_r \to 6q$ and $S_r \to 6\bar{q}$ will then be allowed, providing the source for $B$ asymmetry. These decays occur when the temperature of the universe is $T \sim 0.1 - 100$ GeV. By this time the electroweak sphalerons have gone out of thermal equilibrium, and thus play no role in the $B$ asymmetry generation. We call this mechanism “post–sphaleron baryogenesis”. The three Sakharov conditions for successful baryogenesis [4] are satisfied rather easily in our scheme. The high dimensionality of the $B$ violating coupling of $S_r$ to the quark fields allows the $\Delta B \neq 0$ decays to go out of equilibrium at weak scale temperatures. CP violation occurs in the decay via loop diagrams involving the exchange of the standard model $W^\pm$ gauge bosons. This amplitude has sufficient light quark mass
suppression to explain naturally the observed (small) value of the baryon to photon ratio \( \eta_B \sim 10^{-10} \). The simplest realization of our mechanism involves interactions that violate \( B \) by two units and therefore gives rise to neutron–antineutron oscillations. We find that the successful implementation of our mechanism sets an upper limit on the transition time for \( N \leftrightarrow \bar{N} \) oscillation bringing it to within the realm of observability. This connection provides a strong motivation for improved searches for \( N \leftrightarrow \bar{N} \) oscillation [5].

The connection with \( N \leftrightarrow \bar{N} \) oscillation can be understood as follows. Let us consider an interaction of the form \( S \mathcal{O}_{\Delta B} \), where \( S \) is a standard model singlet complex scalar field (with \( S_r \) denoting its real part) and \( \mathcal{O}_{\Delta B} \) is the baryon number violating operator in question. This interaction will lead to baryon number violation if \( \langle S \rangle \neq 0 \). Suppose the mass dimension of the operator \( S \mathcal{O}_{\Delta B} \) is \( M^{-n} \) with \( n \) positive. The higher the value of \( n \) for an operator \( \mathcal{O}_{\Delta B} \), the lower the mass scale allowed by the existing limits on baryon violation. Since the rate of these \( \Delta B \neq 0 \) interactions in early universe goes like \( M^{-2n} \), the higher the value of \( n \), the easier it is to satisfy the out-of-equilibrium condition at a lower temperature (multi GeV range). Clearly the operator leading to \( B - L \) conserving proton decay mode cannot be useful for us, since present experimental limits on proton lifetime imply that this operator should go out of equilibrium at temperatures of order \( 10^{14} - 10^{15} \) GeV. On the other hand, for a process like \( N \leftrightarrow \bar{N} \) oscillation [6, 7, 8], present experimental lower limits on the oscillation time \( \tau_{N-\bar{N}} \) [9, 10] allow the mass \( M \) appearing in the operator \((fgh)u^c \bar{d}^c d^c f^c \bar{d}^c d^c /M^5\) to be in the multi–TeV range (for the first family Yukawa couplings \( f \sim g \sim h \sim 10^{-3} \)). The out-of-equilibrium temperature for the processes \( S_r \rightarrow 6q \) and \( S_r \rightarrow 6\bar{q} \) is then allowed to be below the sphaleron decoupling temperature of about 100 GeV. We will illustrate how post–sphaleron baryogenesis works using the \( \Delta B = 2 \) process, although the mechanism applies more generally. (The \( \Delta B = 2 \) operator involving left–handed quark doublet fields, \( QQQQQQH^*H^*/M^7 \), has additional Higgs fields and thus a higher dimensionality.)
The high dimensional $\Delta B = 2$ couplings of $S$ are obtained by integrating out colored scalar fields. These colored scalars cannot be much heavier than about a TeV, or else the induced $\eta_B$, consistent with nucleosynthesis limits, will be too small. The prospects for discovering such baryon number carrying colored scalars at the Large Hadron Collider are quite promising.

An attempt to generate baryon asymmetry at a temperature of order MeV via the decay of a heavy ($\sim 50$ TeV) gravitino within supergravity was proposed in Ref. [11]. Such a large gravitino mass would however require fine-tuning to solve the hierarchy problem. Another scenario [12] invokes the decay of the inflaton into squarks, with their subsequent decay producing baryon asymmetry. This mechanism requires that the reheating temperature be less than a GeV in order for the scattering and inverse decays not to wash out the asymmetry. The model presented here differs from these earlier attempts in two crucial ways: (i) There is a strong link between baryon asymmetry and $N \leftrightarrow \bar{N}$ oscillation, and (ii) the mechanism of inducing CP asymmetry via the standard model $W^\pm$ loops which leads naturally to a value of $\eta_B \sim 10^{-10}$ due to light quark mass suppression is entirely new.

2 Light diquarks and observable $N \leftrightarrow \bar{N}$ oscillation

To illustrate our mechanism for post–sphaleron baryogenesis, we consider a generic TeV scale model that gives rise to the higher dimensional operator for $N \leftrightarrow \bar{N}$ oscillation. It consists of the following color sextet, $SU(2)_L$ singlet scalar bosons $(X, Y, Z)$ with hypercharge $-\frac{4}{3}, \frac{2}{3}, \frac{2}{3}$ respectively that couple to the right–handed quarks.\footnote{Color sextet fields are preferred over color triplets, since the sextets do not mediate proton decay.} In addition, there is a complex scalar field $S$ which is a singlet of the standard mode with mass in the 100 GeV range. With this field content one can write down the following standard model invariant
interaction Lagrangian:2
\[ \mathcal{L}_I = \frac{h_{ij}}{2} X d_i^c d_j^c + \frac{f_{ij}}{2} Y u_i^c u_j^c + \frac{g_{ij}}{2} Z (u_i^c d_j^c + u_j^c d_i^c) + \frac{\lambda_1}{2} S X^2 Y + \frac{\lambda_2}{2} S X Z^2 + h.c. \] (1)

If the scalar field $S$ which has $B = 2$ is given a vacuum expectation value, cubic scalar field couplings of the type $X^2 Y$ that break baryon number by two units will be induced. In turn it will lead to $N \leftrightarrow \bar{N}$ oscillation via the diagram of Fig. 1 with $S_r$ replaced by $\langle S \rangle \ [8]$. We note that not all of the $(X,Y,Z)$ fields are needed for $B$ violation and $N \leftrightarrow \bar{N}$ oscillation. $(X,Y)$ or $(X,Z)$ fields will do. We will focus more on these minimal versions in our computation of baryon asymmetry, while for generality we keep all three fields.

To see the constraints on the parameters of the theory, we note that the present limits on $\tau_{N-\bar{N}} \geq 10^8 \text{ sec.}$ implies that the strength $G_{N-\bar{N}}$ of the $\Delta B = 2$ transition is $\leq 10^{-28}$ GeV$^{-5}$. From Fig. 1, we conclude that
\[ G_{N-\bar{N}} \simeq \frac{\lambda_1 \langle S \rangle h_{11}^2 f_{11}}{M_X^2 M_Y^2} + \frac{\lambda_2 \langle S \rangle h_{11} g_{11}^2}{M_X^2 M_Y^2} \leq 10^{-28} \text{ GeV}^{-5}. \] (2)

For $\lambda_{1,2} \sim 1$, $h_{11} \sim f_{11} \sim g_{11} \sim 10^{-3}$, we find $\langle S \rangle \sim M_{X,Y,Z} \simeq 1 \text{ TeV}$ is allowed. In our discussion, we will stay close to this range of parameters and see how one can understand the baryon asymmetry of the universe. In fact, we will see that the masses of $X,Y,Z$ cannot be much larger than a TeV for successful baryogenesis. Note that the couplings $(f,g,h)_{ij}$ to the second and third generation fermions could be larger.

Other constraints can come from low energy observations such as bounds on flavor changing hadronic processes such as $K - \bar{K}$, $D - \bar{D}$ transition etc. This of course depends on any possible mixings between the right handed quark fields, on which we do not have any apriori information. If we make the simplest assumption dictated by the left–right symmetric theories that the left and the right–handed mixings are equal, then the strongest constraints come

An additional term $g_{ij}' Z (u_i^c d_j^c - u_j^c d_i^c)$ with $g_{ij}' = -g_{ji}'$ is also allowed by the standard model gauge symmetry, but this term is forbidden when the model is embedded minimally into a left–right symmetric framework. We do not keep this term explicitly here, its inclusion is however straightforward.
from $K - \bar{K}$ transition which imply that for $h_{11} \sim 10^{-3}$, $M_X \geq 1$ TeV, which is consistent with our choice of parameters dictated by observability of $N \leftrightarrow \bar{N}$ transition.

The model of Eq. (1) is embeddable into an $SU(2)_L \times SU(2)_R \times SU(4)_c$ framework where the quarks and leptons transform as $\psi : (2, 1, 4) \oplus \psi^c : (1, 2, \bar{4})$ representations and the Higgs fields $X, Y, Z, S$ are part of the $\Delta^c : (1, 3, 10)$ multiplet. In fact the $S$ field corresponds to the $\Delta^c_{\nu\nu}$ component that acquires a vev and breaks $B - L$ by two units. In this paper, we will not discuss the full set of constraints that arise in this embedding but rather simply work within the scalar field model described in Eq. (1). All our conclusions below apply to the $SU(2)_L \times SU(2)_R \times SU(4)_c$ model as well. While we take baryon number as part of the gauge symmetry, the mechanism of $B$ asymmetry generation also works if $B$ is a spontaneously broken global symmetry as in Ref. [13].

3 Origin of matter

Before proceeding to the discussion of how baryon asymmetry arises in this model, let us first consider the effect of the new interactions in Eq. (1) on any pre-existing baryon asymmetry. For this purpose, we assume the following mass hierarchy between the $S$ field and the $(X, Y, Z)$ fields: $M_S \sim 100$ GeV $\ll M_{X,Y,Z} \sim$ TeV. For $T \geq M_{X,Y,Z}$, the $\Delta B = 2$ interaction rates scale like $T$ and are in equilibrium at least down to $T \simeq M_{X,Y,Z}$. They will therefore erase any pre-existing baryon asymmetry. They remain in equilibrium down to the temperature $T_*$ determined by the inequality:

$$\frac{1}{(2\pi)^9} \frac{\lambda_{ij}^2 h_{ij}^2 g_{ij}^4 T_{13}^3}{M_{X,Y,Z}^2} \leq \frac{g_*^{1/2} T^2}{M_{Pl}}.\tag{3}$$

Here $h$ and $g$ refer to the largest of $h_{ij}$ or $g_{ij}$ ($i, j$ are family indices). For $h, g$ in the range of $0.1 - 1$, this leads to $T_* \simeq (0.6 - 0.2) M_{X,Y,Z}$.

The singlet field $S$ will play a key role in the generation of baryon asymmetry. We assume
that \(\langle S \rangle \sim M_X\) and \(M_{S_r} \sim 10^2\) GeV, where \(S_r\) is the real part of the \(S\) field after its vev is subtracted. \(S_r\) can then decay into final states with \(B = \pm 2\), viz., \(S_r \rightarrow 6q\) and \(S_r \rightarrow 6\bar{q}\), inducing a net baryon asymmetry.

On the way to calculating the baryon asymmetry, let us first discuss the out of equilibrium condition. As the temperature of the universe falls below the masses of the \(X, Y, Z\) particles, the annihilation processes \(X\bar{X} \rightarrow d\bar{c}d\bar{c}\) (and analogous processes for \(Y\) and \(Z\)) remain in equilibrium. As a result, the number density of \(X, Y, Z\) particles gets depleted and only the \(S\) particle survives along with the usual standard model particles. The primary decay modes of \(S_r\) are \(S_r \rightarrow u^cd\bar{c}d\bar{c}\) and \(S_r \rightarrow \bar{u}^c\bar{d}\bar{c}\bar{d}\bar{c}\). There could be other decay modes which depend on the details of the model. Those can be made negligible by choice of parameters which do not affect our discussions of \(N \leftrightarrow \bar{N}\) oscillation and baryogenesis. We will discuss them later in the paper. For \(T \geq M_{S_r}\), the decay rate of \(S_r\) is given by the left–hand side of Eq. (3). This decay goes out of equilibrium around \(T_\star \sim 0.4M_X\), or around 500 GeV. Below this temperature the decay rate of \(S_r\) falls very rapidly as the temperature cools. However as soon as \(T \leq M_{S_r}\), the decay rate becomes a constant while the expansion rate of the universe slows down. So at a temperature \(T_d\), \(S\) will start to decay where \(T_d\) is given by

\[
T_d \approx \left[ \frac{18P\lambda_2^2h^2g^4 M_{Pl} M_{S_r}^{13}}{(2\pi)^91.66g_s^{1/2}(6M_X)^{12}} \right]^{1/2}.
\]

This is obtained by equating the decay rate of \(S_r\) to the expansion rate of the universe. In Eq. (4) the factor 18 is a color factor, \(h^2 = \text{Tr}(h^\dagger h)\), etc, while \(P\) is a phase space factor, which we have computed for the six body decay via Monte Carlo methods and found \(P \approx 2.05\). The corresponding epoch must be above that of big bang nucleosynthesis. This puts a constraint on the parameters of the model. For instance, for \(M_S \sim 200\) GeV and \(M_X \sim \text{TeV}\), we get \(T_d \sim 40\) MeV (for \(g \sim h \sim 1\)). We will conduct the rest of the discussion with this set of parameters as a representative set. Note that \(M_X\) cannot be much larger than about 1
Figure 1: Tree level diagrams contributing to $S_r$ decays into 6 anti-quarks. There are other diagrams where $S_r$ decays into 6 quarks, obtained from the above by reversing the arrows of the quark fields.

TeV, otherwise $T_d$ will be below few MeV, affecting big bang nucleosynthesis significantly. Note also that the at least some of the couplings in $h$ and $g$ should be of order one. This would imply that the first family couplings should be of order $(10^{-3} - 10^{-4})$ from naturalness ($h_{11} \sim V_{td}^2 h_{33}$ etc), making $N \leftrightarrow \overline{N}$ oscillation accessible to next generation experiments.

We now proceed to calculate the baryon asymmetry in this model. It is well known that baryon asymmetry can arise only via the interference of a tree diagram with a one loop diagram which has an absorptive part. The tree diagrams are clearly the one where $S_r \rightarrow 6\bar{q}$ and $S_r \rightarrow 6q$. There are however two classes of loop diagrams that can contribute to baryon asymmetry: one where the loop involves the same fields $X, Y$ and $Z$ as in Fig. 2 (there is a second loop diagram of this type with $(X, Z)$ fields inside the loop), and a second one involving $W^{\pm}$ gauge boson exchange as shown in Fig. 3. In the $(X, Z)$ model and in the $(X, Y)$ model, only the latter contribution exists (the former trace being real). So we focus on that latter, which involves only standard model physics at this scale and has the advantage that it involves less number of arbitrary parameters.
We summarize the results of our calculations for the $W^\pm$ exchange diagrams. If one of the external up–type quarks is the top quark, the corresponding quark line receives a wave function correction via $W^\pm$ gauge boson exchange. The baryon asymmetry from this diagram is found to be

$$\epsilon_B^{\text{wave}} \simeq -\frac{3\alpha_2}{8} \left( 1 + \frac{m_W^4}{m_t^4} \right) \frac{\text{Im} \left[ V^* \hat{M}_u^2 V^T \hat{M}_u g g^\dagger \right]}{m_t m_W^2 (gg^\dagger)_{33}}$$

where $\hat{M}_u = \text{diag}(m_u, m_c, m_t)$, $\hat{M}_d = \text{diag}(m_d, m_s, m_b)$ and $V$ is the CKM matrix. $\text{Br}$ stands for the branching ratio of $S_\tau$ into $6q + \bar{6}q$.

The vertex correction via the $W$ boson exchange gives an asymmetry given by

$$\epsilon_B^{\text{vertex}} \simeq -\frac{\alpha_2 \text{Im} \text{Tr}[g^T \hat{M}_u V g^\dagger V^* \hat{M}_d]}{4 \text{Tr}(g^T g)}.$$  \hspace{1cm} (6)

Here we have assumed that $M_{S_\tau} \gg m_t$. In the limit where $m_{S_\tau} \ll m_W$, we have the same asymmetry as in Eq. (6), but with a factor of $(-1/4)$ multiplying it. Of course in this case, decays involving final state top quark are disallowed, which is to be implemented by removing the top quark contribution in the trace of Eq. (6).\(^3\)

\(^3\)These $W^\pm$ loops do not conflict with the theorem of Ref. [14] which states that no baryon asymmetry
It is interesting to note that in this mechanism, there is a natural explanation of the observed baryon asymmetry \( \eta_B \sim 10^{-10} \). It follows from the light quark mass and mixing angle suppression. As an example, consider the following choice of parameters: \( m_c(m_c) = 1.27 \text{ GeV}, \ m_b(m_b) = 4.25 \text{ GeV}, \ m_t = 174 \text{ GeV}, \ V_{cb} \simeq 0.04, \ M_S = 200 \text{ GeV} \) and \( |g_{33}| \simeq |g_{23}| \sim 1 \), with smaller values of \( g_{1i} \). We find \( \varepsilon_B \sim 10^{-8} \) in this case from Eq. (5). The corresponding value from Eq. (6) is an order of magnitude larger, for the same input parameters.

There is a further dilution of the baryon asymmetry arising from the fact that \( T_d \ll M_{S_r} \) since the decay of \( S_r \) also releases entropy into the universe. In this case the baryon asymmetry reads

\[
\eta_B \simeq \frac{T_d}{M_{S_r}} \epsilon_B.
\]  

(7)

In order that this dilution effect is not excessive, there must be a lower limit on the ratio \( T_d/M_{S_r} \). From our estimate above we require that \( T_d/M_{S_r} \geq 0.01 \). Since the decay rate of the \( S_r \) boson depends inversely as a high power of \( M_{X,Y} \), higher \( X,Y \) bosons would imply that \( \Gamma_S \sim H \) is satisfied at a lower temperature and hence give a lower \( T_d/M_{S_r} \). In Fig. 4 can be induced by dressing a \( \Delta B = 1 \) vertex by baryon number conserving interactions. Since \( S_r \) field has no definite baryon number, owing to \( \langle S \rangle \neq 0 \), the theorem is not applicable in our case.
Figure 4: The allowed range of $M_X$ and $M_S$ needed to generate the baryon asymmetry (along the black curve), decay temperature above 200 $MeV$ (points below the dashed curve) and $\tau_{N\bar{N}} \geq 10^8$ sec with $\bar{\lambda} = 10^{-4}$ (points above the red curve).

we have plotted $M_{X,Y}$ vs $M_{S_\tau}$ which gives the right amount of baryon asymmetry which is consistent with the demand that the decay of $S_\tau$ occurs before the QCD phase transition (i.e $T_d \geq 200$ $MeV$). Using the effective coupling $\bar{\lambda} \equiv (\lambda_1 h^2_{11} f_{11})^{1/4} \sim (\lambda_2 h g^2_{11})^{1/4}$ to be of order $10^{-4}$ implies that the $10^9$ sec $\leq \tau_{N-\bar{N}} \leq 10^{11}$ sec for $M_{S_\tau}$ in range of $\simeq 100 - 300$ GeV.

We now comment briefly on some other aspects of the model:

(i) If the $(X,Y,Z)$ scalars are all present, the loop diagram of Fig. 2 will contribute to baryon asymmetry. Since there are two diagrams of the type shown in Fig. 2 the relevant trace has an imaginary piece. The asymmetry will have a suppression factor $M_{S_\tau}^2/M_X^2$, in addition to a loop factor and the Yukawa suppression. Although not very predictive, this model can yield adequate baryon asymmetry.

(ii) The singlet field $S$ can have a renormalizable interaction with the standard model Higgs doublet of the form $\lambda_S S^\dagger S H^\dagger H$. After the fields $S$ and $H$ acquire vacuum expectation
values, the Re $S$ and the SM Higgs field can mix with each other opening up new decay channels for the $S_r$ field such as $S_r \rightarrow b\bar{b}$ etc. We estimate this decay width to be

$$\Gamma(S_r \rightarrow b\bar{b}) \sim \frac{3\beta^2 m_b^2 M_S}{4\pi M_W^2}$$

where $\beta$ is the $S_r - h$ mixing angle. There are two constraints on this decay mode in order for our scenario to work: first, this could contribute dominantly to the $S_r$ decay width thereby diluting the baryon asymmetry. Second, this model should be out of equilibrium at $T = M_S$. If the model is non-supersymmetric, these two conditions are satisfied if $\lambda_S \leq 10^{-6}$. This coupling is automatically forbidden if the model is supersymmetric.

(iii) The present considerations could be easily extended to include supersymmetry as part of the new physics beyond the standard model. The $SX^2Y$ and $SXZZ$ interactions in this case are nonrenormalizable [15]. However, in this case we also expect mass terms in the superpotential of the form $M_S S \bar{S}$ so that the effective four scalar interaction responsible for baryogenesis is in the same form as discussed above.

(iv) Our theory is also testable in collider experiments such as the LHC since we have colored diquark scalar fields with masses in the TeV range. In a $p\bar{p}$ collision one could produce the $X, Y, Z$ bosons either in pairs via the process $q\bar{q} \rightarrow X\bar{X}$ or singly via the process $q + g \rightarrow X + \bar{q}$. In the first case the signal would be a four jet final state whereas in the second case, it would be three jet final states. It would therefore be important to search for such final states at LHC. One distinguishing feature of these bosons is that they carry baryon number, which may be testable in the decays of these bosons into top quark and bottom quarks.
4 Conclusion

In conclusion, we have presented a new mechanism for baryogenesis which does not rely on the electroweak sphalerons but rather directly produces matter–antimatter asymmetry using higher dimensional baryon violating couplings of a scalar field. The baryon asymmetry is produced at the weak scale. This mechanism can be tested by searches for baryon violating processes such as neutron antineutron oscillation, as well as by the discovery of colored scalars at the LHC.

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