ELECTROGENESIS IN A SCALAR FIELD DOMINATED EPOCH

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Abstract. In this talk I discuss models in which a homogeneous scalar field is used to modify standard cosmology above the nucleosynthesis scale to provide an explanation for the observed matter-antimatter asymmetry of the Universe.

I INTRODUCTION

Scalar fields are used to model either a very early universe (inflation), or a very late universe (‘quintessence’) [1]. Their role at intermediate times is however largely ignored. In this talk I discuss models in which a homogeneous scalar field is used to modify the standard cosmology at the electroweak scale to allow for baryogenesis, even when the electroweak transition is smooth or weakly first order.

In order to produce any baryon number, a source is required that drives the Universe out of equilibrium [5]. Since at the electroweak scale the expansion rate is very small ($H/T \sim 10^{-16}$), it is often assumed that it cannot drive baryogenesis, simply because the produced baryon-to-entropy ratio, $n_B/s \propto H/T$, is too small to account for observation, $(n_B/s)_{observed} \sim 5 \times 10^{-11}$.

II SCALAR FIELDS AND THE EXPANSION RATE

In Refs. [2] we discussed models in which, based on the dominance of a kinetic scalar field mode (kination), the expansion rate of the Universe changes to

\[ \frac{H}{T} \sim \frac{T}{T_{\text{reh}}} \left( \frac{H}{T} \right)_{\text{rad}}, \]

where $(H/T)_{\text{rad}}$ is the expansion rate in radiation-dominated universe, and $T_{\text{reh}}$ is the ‘reheat’ temperature at which the energy-densities are equal, $\rho_\phi \sim \rho_{\text{rad}}$ (see

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1) Invited talks presented at the conferences Strong and Electroweak Matter (SEWM-2000) in Marseille, (June 14-17, 2000), and Cosmology and Astroparticle Physics (CAPP-2000) in Verbier, Switzerland (July 17-28, 2000); based on work with Michael Joyce [2–4].
Model A in figure 1). At the nucleosynthesis scale, $T_{ns} \sim 1\text{MeV}$, the Universe is radiation dominated, and hence $T_{\text{reh}} \geq T_{ns}$. Since $T_{ew}/T_{ns} \sim 10^5$, the expansion rate at the electroweak scale ($T_{ew} \sim 100\text{GeV}$) can be enhanced to about $(H/T)_{ew} \sim 10^{-11}$. With some tuning in the parameters of the model [2], this is enough for successful baryogenesis even at a smooth or a weakly first order transition. We also note that the same scalar field can be both the inflaton and the kinaton [2].

![Figure 1](image-url)

**FIGURE 1.** Evolution of energy density in radiation and the dominant scalar field as a function of temperature. Two cases are illustrated: Model (A) in which the dominant scalar component scales faster then radiation, but does not decay (solid lines), and Model (B) in which the scalar field decays (dashed lines).

The expansion rate can be further enhanced if the scalar field decays. In this case we have

$$
\frac{H}{T} \sim \left( \frac{T}{T_{\text{decay}}} \right)^2 \left( \frac{H}{T} \right)_{\text{rad}},
$$

where again $T_{ew} \geq T_{\text{decay}} \geq T_{ns}$, and $T_{\text{decay}} = T_{\text{reh}}$ is the temperature at which the field decays (see Model B in figure 1). At a first sight one can get the expansion rate at the electroweak scale high enough to drive baryogenesis. There is a caveat however: as the scalar field decays, entropy is released, which then dilutes the original baryon number produced at the electroweak scale. The entropy release is minimal if the scalar field energy is dominated by the kinetic mode [3].

**III APPROXIMATELY CONSERVED CHARGES AND BARYON NUMBER**

We have argued that the dominance of a kinetic scalar mode can be used to increase the expansion rate of the Universe at the electroweak scale by orders of
magnitude. In fact the expansion rate can easily become larger than the destruction rate of some of approximately conserved species, e.g.

\[ H(T_{ew}) \geq \Gamma_{e_R}, \Gamma_{\mu_R}, \Gamma_{u_R}, ... \]  

(3)

This simply means that, if any of these charges are produced above the electroweak scale, they decay only below the electroweak scale, that is when the baryon number violating processes are already frozen-in. If a net right-handed electron number, \( e_R \), is produced at a scale \( T > T_{ew} \), in chemical equilibrium the baryon number \( B \) is shifted to [4]

\[ B \approx \frac{1}{3} e_R. \]  

(4)

This simple relation describes the correct local chemical equilibrium as long as the destruction rate for the right-handed electrons is large compared with the expansion rate at the electroweak scale, i.e. \( \Gamma_{e_R} \sim 10^{-13}T_{ew} < H(T_{ew}) \). At \( T = T_{ew} \) the baryon-number violating processes (‘sphalerons’) freeze-in, i.e. the sphaleron rate drops below the expansion rate, and the baryon number (4) remains frozen until today. Eq. (4) can be intuitively understood as follows. Above the electroweak scale the Universe must be hypercharge-neutral, \( Y = 0 \). In the presence of net \( e_R \) the corresponding electron hypercharge \( Y_{e_R} = y_{e_R}e_R \) must be screened by various charges pulled out of the plasma. The charges that minimize the relevant free energy include the quarks that carry a net baryon number \( B \) as given in Eq. (4).

**IV ELECTROGENESIS**

We now present a simple perturbative model for production of the right-handed electrons required for baryogenesis (cf. Eq. (4)). To this purpose we introduce heavy scalar fields \( \Phi^a \) (with a mass at least in the TeV range) that couple to the standard-model fermions via a Yukawa interaction term of the form

\[ \mathcal{L}_{CP}[\Phi^a] = -h_{ij}^a \Phi^a \bar{\psi}_i \psi_j + h.c., \]  

(5)

where the couplings \( h_{ij}^a \) are CP violating, i.e. \( h_{ij}^a \neq h_{ji}^a \) (\( h^a \) denotes the matrix of couplings). In order to violate CP symmetry, \( h^a \) must contain a complex phase unremovable by phase transformations on the whole Lagrangian, which can be achieved by the flavor mixing structure and the existence of at least two such scalars. The most stringent constraints on the masses and the couplings of such scalars come from the fact that they are flavor changing. For leptons the strongest constraint of this type comes from the bounds on the decay \( \mu \rightarrow e\gamma \). For couplings \( h_{ij}^a \) of order one this requires masses \( M_\Phi \geq 10^6 \text{TeV} \).

When \( \Phi^a \) decay out of equilibrium, a net \( e_R \) may be produced. An example of such a decay channel is shown in figure 2, where CP violation is realised as the
FIGURE 2. Tree and one loop diagrams for the three body decay $\Phi \rightarrow \bar{e}_R\Psi_R^+\Phi'$, with the appropriate couplings at the vertices. We assume that $\Phi$ is heavier than $\Phi'$. When the second outgoing lepton is a $\mu$ or $\tau$ lepton the process produces net $e_R$ number.

interference term between the tree level and 1-loop 3-body decay channels (cf. [6]). The resulting electron-to-entropy ratio is then

$$\frac{e_R}{s} \sim \frac{10^{-2}}{g_*} |h|^4 \delta_{CP},$$

where $g_*$ is the number of relativistic degrees of freedom in the plasma, $\delta_{CP}$ the relevant CP-violating angle. In order that $\Phi^a$ decay out of equilibrium, they ought to be sufficiently massive. One finds [4] that in Model A,

$$M_{\Phi} > 5 |h| \times 10^6 \text{ GeV},$$

while in Model B, when the dominant component decays, the bound reads

$$M_{\Phi} > 3 |h|^2 \text{ TeV}.$$ 

This implies that the scalars $\Phi^a$ may be observable by the future accelerators (e.g. LHC). This fact alone gives a sufficient motivation for a more detailed investigation of the models that contain such heavy scalar fields.

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