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Abstract

I report on a critical analysis of the scenario of electroweak baryogenesis mediated by nonsuperconducting cosmic strings. This mechanism relies upon electroweak symmetry restoration in a region around cosmic strings, where sphalerons would be unsuppressed. I discuss the various problems this scenario has to face, presenting a careful computation of the sphaleron rates inside the strings, of the chemical potential for chiral number and of the efficiency of baryogenesis in different regimes of string networks. The conclusion is that the asymmetry in baryon number generated by this scenario is smaller than the observed value by at least 10 orders of magnitude.

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I report on a critical analysis of the scenario of electroweak baryogenesis mediated by nonsuperconducting cosmic strings. This mechanism relies upon electroweak symmetry restoration in a region around cosmic strings, where sphalerons would be unsuppressed. I discuss the various problems this scenario has to face, presenting a careful computation of the sphaleron rates inside the strings, of the chemical potential for chiral number and of the efficiency of baryogenesis in different regimes of string networks. The conclusion is that the asymmetry in baryon number generated by this scenario is smaller than the observed value by at least 10 orders of magnitude.

1 Introduction

Electroweak baryogenesis [1] is a beautiful idea which fails (or is about to fail) in the best motivated models we have for physics at the Fermi scale ($\sim 100\, GeV$). In the Standard Model, LEP II experiments set a lower bound on the mass of the Higgs boson of about 97 GeV, implying that the electroweak phase transition in that model is not first order but rather a crossover [2]. In the Minimal Supersymmetric Standard Model the electroweak phase transition can be first order and sufficiently strong to allow for electroweak baryogenesis, but this occurs in a very small region of parameter space [3] which presumably will be ruled out by LEP II in a couple of years.

One may take the previous negative results as indication that the asymmetry in baryon number was not created at the electroweak epoch, but rather related to the physics of $B-L$ violation and neutrino masses. To stick to electroweak baryogenesis one can consider extensions of the particle content of the model to get a stronger electroweak phase transition (e.g. extensions which include singlets). In this talk I will consider another possibility: how the remnants of physics at energy scales higher than the electroweak scale (cosmic strings in this case) can be useful to overcome the problems of having a weak electroweak phase transition.

Electroweak baryogenesis requires the co-existence of regions of large and small $\langle\varphi\rangle/T$, where T is the temperature and $\langle\varphi\rangle$ the (T-dependent) Higgs vacuum expectation value. At small or zero $\langle\varphi\rangle/T$ sphalerons are unsuppressed and mediate baryon number violation, while large $\langle\varphi\rangle/T$ is needed to store the created baryon number (for $\langle\varphi\rangle/T \geq 1$ sphaleron transitions are ineffective and baryon number is conserved). Below the critical temperature T_c^{EW} of the electroweak phase transition and irrespective of whether it is first or second order, $\langle\varphi\rangle/T$ grows until sphaleron transitions are shut-off. For baryogenesis to be possible at those times, we need some region where $\langle\varphi\rangle$ is forced to remain zero or small. The idea we examine in this talk is that this can be the case along topological defects (like cosmic strings) left over from some other cosmological phase transition that took place before the electroweak epoch [4]. If the electroweak symmetry is restored in some region around the strings, sphalerons could be unsuppressed in the string cores while they would be ineffective in the bulk of space, away from the strings. The motion of the string network, in a similar way as the motion of bubble walls in the usual first-order phase-transition scenario, will leave a trail of net baryon number behind.

Some problems with this scenario come immediately to mind. First, it is clear that the space swept by the defects is much smaller than the total volume, so there will be a geometrical suppression factor with respect to the usual bubble-mediated scenario [4]. Another suppression factor arises from the fact that there is a partial cancellation between front and back walls of the string, which tend to produce asymmetries of opposite signs [4]. Another problem comes from the condition that the symmetry restoration region (which naively would be of size $R_{rest} \sim 1/\sqrt{\lambda}\langle\varphi\rangle$, where λ is the quartic Higgs coupling) should be large enough to contain sphalerons (which in the symmetric phase have size $R_{sph} \sim 1/g^2T$), while outside the strings, sphalerons should be suppressed ($\langle\varphi\rangle/T \geq 1$). Combining both conditions one obtains $\lambda \leq g^4$, which means the scenario would require small values of the Higgs mass, in conflict with experimental bounds. LEP II tells us that λ is at least of order g^2, so that sphalerons won't fit in the restoration region. In other words, for realistic values of the Higgs mass sphalerons are not going to be fully unsuppressed. We will measure how effective they are by writing the rate of sphaleron transitions per unit time and unit of string length as $\Gamma_l = \kappa_l \alpha_w^2 T^2$. For a string with $R_{rest} = R_{sph}$, one has $\Gamma_l R_{rest}^2$ equal to the rate in the symmetric phase, corresponding to $\kappa_l \sim 1$. Values of κ_l much smaller than 1 would mean that sphalerons are not really unsuppressed inside the strings.

In the rest of the talk I review the careful analysis of this mechanism contained in ref. [5], to which I refer the interested reader for further details.

2 Strings with electroweak symmetry restoration

Cosmic strings [6] are 1-dimensional solitons, stable by topological reasons, that can form in the spontaneous breaking of a symmetry G where I consider the simplest case, $G = U(1)$, in this talk. A model with a complex scalar S and lagrangian

$$\mathcal{L} = \partial_\mu S^* \partial^\mu S - \lambda_S (S^* S - S_0^2)^2, \tag{1}$$

admits global strings: configurations with $S = 0$ along some line (say the z-axis) and $S(r) = f(r) S_0 e^{i\theta}$, with $f(\infty) \to 1$, where r is the distance to the z-axis and θ the azimuthal angle. The radius of these strings (where most of the energy is trapped) is set by the scale $1/m_S \equiv 1/\sqrt{\lambda_S} S_0$.

If the $U(1)$ is made local, in addition to the S field, a non-zero gauge field is also present, $A_\mu = -a(r)\partial_\mu \theta / q_S$, with $a(\infty) = 1$, where q_S is the $U(1)$ charge of the S field. This gauge field is such that the covariant derivative $D_\mu S$ goes to zero for large r resulting in a finite energy per unit length of string.

We assume that S-strings (global or local) form at some temperature $T_c^S > T_c^{EW}$ and are present at the time of the electroweak phase transition. To force $\langle \varphi \rangle \to 0$ in the cores of the strings, the Higgs field must interact either with the S field or the A_μ field (if the strings are local):

2.1 $S - \varphi$ interaction

Suppose the scalar potential has the form

$$V(S,\varphi) = \lambda_S (|S|^2 - S_0^2)^2 - \gamma (|S|^2 - S_0^2)(|\varphi|^2 - \varphi_0^2) + \lambda (|\varphi|^2 - \varphi_0^2)^2, \tag{2}$$

with $\gamma > 0$. The mass squared of the Higgs field in the string background is $m_\varphi^2(r) \sim \gamma (S_0^2 - |S(r)|^2) - 2\lambda \varphi_0^2$, which is negative outside the string core but can be positive inside, so that electroweak symmetry tends to be restored along the strings. Exploring the $(S_0, \lambda_S, \gamma, \lambda)$ parameter space, the typical case, with $\lambda_S S_0^2 \gg \lambda \varphi_0^2$ leads to $R_{rest} \sim 1/m_\varphi(\infty)$. The best posible case to get a large restoration region has $\lambda_S \ll \gamma \ll \lambda$ and $S_0 \gg \varphi_0$ and gives $R_{rest} \sim \sqrt{\gamma/\lambda_S}/m_\varphi(\infty)$.

2.2 $S - A_\mu$ interaction

In this case we assume that the Higgs field carries a charge q_φ under the extra $U(1)$ responsible for the strings, so that its covariant derivative has an extra piece. As we saw, the A_μ field in the string goes like $-1/q_S r$ at large r to cancel the azimuthal derivative of S, give vanishing $D_\mu S$ and minimize energy. In $D_\mu \varphi$, the A_μ contribution is now proportional to q_φ / q_S and the azimuthal derivative of φ can cancel $D_\mu \varphi$ only if q_φ/q_S is an integer. If that is not the case, a Z_μ boson condensate is induced until the covariant derivative is cancelled [7]. In any case, a non-zero winding of φ forces $\varphi \to 0$ in the string core ($r = 0$). The restoration region around $r = 0$ is larger in the presence of a non-zero Z_μ string (case of non-integer q_φ/q_S).

3 Sphaleron rates and CP asymmetry in the string cores

In general, with no tuning of potential parameters nor a Z_μ condensate, $\langle\varphi\rangle$ is zero only at the string core ($r = 0$) and rises inmediately away from that line. As the symmetry is never really restored in a wide region, the energy of the sphaleron in such background (it can be computed in the lattice looking for a saddle point of the energy functional) is only about a factor 0.7 smaller than the sphaleron energy in the broken phase (alternatively $\kappa_l \sim 10^{-6}$: that is, sphalerons are not really unsuppressed in this type of strings).

The situation is better when a Z_μ-field is induced, in which case $\kappa_l \sim 1/30$ for $\langle\varphi\rangle/T \sim 1$ (this number can be obtained in the lattice using a fully non-perturbative approach and tracking Chern-Simons number in real time evolution). However this number is very sensitive to T and drops significantly when T decreases.

Fully unsuppressed sphalerons can only be obtained in the global $U(1)$ case for large enough γ/λ_S. In fact, to obtain an asymmetry of the order of the observed one, one would need $\gamma/\lambda_S \sim 10^{14}$. On the other hand, stability of the potential requires $4\lambda/\gamma > \gamma/\lambda_S$, so that $\lambda/\lambda_S \sim 10^{28}$. Such an ad-hoc and wild fine-tuning of the parameters prevents us from taking this particular case seriously.

Unsuppressed sphaleron transitions inside the string cores are not sufficient to generate the baryon asymmetry: they must occur in a background with CP asymmetric particle distributions so that the sign of the B-violation is biased. This asymmetry comes about if the interactions between the particles in the plasma and the string walls violate CP. In that case the walls of a moving string act as sources of chiral-number flux (which would be zero if the string velocity v_S were zero). This asymmetry diffuses away from the walls and only that inside the string is useful to create baryons (for geometrical reasons it is also clear that this diffusion effect is less efficient for strings than for bubbles). In conclusion, we have to compute the chemical potential μ for chiral number inside the strings. General arguments (confirmed by detailed analysis of particular models) give the result $\mu = K v_S^2 T$ for small v_S, with $K \lesssim 0.01$ and $\mu = K' T$ for $v_S \sim 1$ with K' of order 1.

4 Evolution of string networks and efficiency of baryogenesis

To get a final number for the asymmetry generated by this mechanism, we need to know how many strings there are and how quickly they are moving (the best case being that of a dense network of fast moving strings). We can describe the string network by a mean average separation between strings $R(t)$ and a mean average velocity $v_S(t)$. The evolution of these quantities with time t is governed by Hubble expansion ($H \sim 1/2t$); energy loss by loop formation; and friction with the plasma. The friction force goes like $F \sim v_S T^3$: it is important at early times when it dominates the dynamics of the evolution. This is the friction dominated or Kibble regime, with $R(t) \sim t^{5/4}$ and $v_S(t) \sim t^{1/4} \sim HR(t)$. Eventually, friction will no longer be important and a scaling regime is reached with $R(t) \sim 1/H$ and $v_S \sim 1$.

In conclusion, to get the final number for the baryon asymmetry we start with the equation for the rate of change of baryon number N_B per unit time and unit length of string:

$$\frac{dN_B}{dLdt} = 1.5[\kappa_l \alpha_w^2 T^2]\frac{\mu}{T}. \tag{3}$$

If we use the results for κ_l and μ previously discussed, and integrate eq.(3) in one Hubble time (this is because κ_l is shut-off quickly with decreasing T) using the network evolution results just presented we end up with the result that

$$\left[\frac{N_B}{N_\gamma}\right]_{strings} \lesssim 10^{-10}\left[\frac{N_B}{N_\gamma}\right]_{observed}. \tag{4}$$

That is, the mechanism just studied is uncapable of generating a sufficiently large matter-antimatter asymmetry.

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