

A New Mechanism for Baryogenesis in Low Energy Supersymmetry Breaking Models

R. Brandenberger^{1*} and A. Riotto^{2†}

¹*Physics Department, Brown University, Providence, RI. 02912, USA.*

²*Theory Group, CERN, CH-1211 Genève 23, Switzerland.*

Abstract

A generic prediction of models where supersymmetry is broken at scales within a few orders of magnitude of the weak scale and is fed down to the observable sector by gauge interactions is the existence of superconducting cosmic strings which carry baryon number. In this paper we propose a novel mechanism for the generation of the baryon asymmetry which takes place at temperatures much lower than the weak scale. Superconducting strings act like “bags” containing the baryon charge and protect it from sphaleron wash-out throughout the evolution of the Universe, until baryon number violating processes become harmless. This mechanism is efficient even if the electroweak phase transition in the MSSM is of the second order and therefore does not impose any upper bound on the mass of the Higgs boson.

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*rhb@het.brown.edu.

†riotto@nxth04.cern.ch, on leave of absence from Theoretical Physics Dept., Oxford Univ., Oxford, U.K.

I. INTRODUCTION

Baryogenesis is an important channel which can relate particle physics beyond the standard model to data. The challenge for particle physics is to determine how the observed baryon to entropy ratio of $n_B/s \sim 10^{-10}$ can be generated from symmetric initial conditions ($n_B = 0$) in the very early Universe (here, n_B and s denote the net baryon number and entropy density, respectively).

For baryogenesis to be possible, it is necessary for the particle physics model to admit baryon number violating processes, to have C and CP violation, and for the relevant processes to occur out of thermal equilibrium. If $B - L$ is a symmetry of the theory, then [1] sphaleron processes which are unsuppressed above the electroweak symmetry breaking scale η_{EW} will erase any baryon asymmetry generated at temperatures higher than T_{EW} , the temperature corresponding to η_{EW} , and the presently observed n_B/s ratio must be generated below T_{EW} (for a recent review of baryogenesis, see Ref. [2]).

A lot of recent work (see e.g. the review articles [3–5]) has focused on ways of generating the baryon asymmetry by electroweak sphaleron processes at or below T_{EW} . It is almost universally agreed (see, however, Ref. [6] for a dissenting view) that this is not possible in the standard model, firstly because CP violation couples to sphaleron processes with insufficient strength, secondly because the electroweak phase transition is not strongly first order and hence does not proceed by the nucleation of critical bubbles. Such critical bubbles are the sites of non-equilibrium which are usually used for electroweak baryogenesis. The constraints of demanding that the phase transition be strongly first order, and that baryon nonconserving processes be out of equilibrium immediately below T_{EW} are hard to reconcile even in extensions of the standard model. For instance, a successful electroweak baryogenesis within the Minimal Supersymmetric Standard Model (MSSM) requires the Higgs boson and the right-handed stop to be light and CP-violating phases of the order of 0.1 [7].

There is, however, another mechanism which can be effective in mediating baryogenesis. Topological defects are out-of-equilibrium field configurations which can trigger baryogenesis

in several ways: they can act as field condensates which release baryons during their decay [8], or they can be used as the catalysis sites of sphaleron processes in a way similar to how bubble walls function [9–11].

In this paper, we work out an idea first suggested in [12] that defects may play a crucial role in baryogenesis in a class of supersymmetric theories where supersymmetry breaking occurs at a low energy scale. It has been shown in previous work [12] that these models admit superconducting cosmic strings with baryon number carrying condensates.

As we will discuss, these strings are able to trap the baryon number from the time of string formation ($T = T_c$) until the temperature of the Universe has dropped down to values much lower than T_{EW} , when sphaleron processes have fallen out of equilibrium. The generation of the baryon asymmetry takes place when the baryon number carried by the condensate in the core of the strings is released in the thermal bath at very low temperatures (this mechanism is different from defect-mediated electroweak baryogenesis [9–11], another mechanism which can be implemented in certain supersymmetric models [13]). The out-of-equilibrium condition is therefore naturally attained. A significant advantage of defect-mediated baryogenesis is that the condition that sphaleron processes must be out of equilibrium immediately below T_{EW} is no longer required [8] and therefore the stringent conditions on the MSSM Higgs spectrum may be relaxed.

Supersymmetric extensions are the best-motivated of the particle physics theories beyond the standard model. Supersymmetry solves the gauge hierarchy problem, couples gauge theory to gravity, and may easily generate the gauge coupling unification. However, the present state of the Universe is not supersymmetric, and thus supersymmetry must be broken in Nature at some scale higher than $\sim \text{TeV}$.

Spontaneous supersymmetry breaking entails adding to the standard model and its supersymmetric copy some new fields acting as a supersymmetry breaking sector. The supersymmetry breaking sector cannot have any renormalizable couplings to the “visible” sector, and is therefore often called “hidden” or “secluded”. There are various ways in which supersymmetry breaking can be communicated from the “hidden” to the “visible” sector (see e.g.

[14] for a recent review): supergravity mediation, gauge mediation, and mediation via $U(1)$ gauge factors. Since in the most general and natural scenarios of low energy supersymmetry breaking there are two different sources of contributions to ordinary superpartner masses, gauge mediation from the messenger sector and anomalous $U(1)$ D -terms, a novel feature is the prediction of superconducting cosmic strings [12]. The typical energy scale of these strings is about 10^2 TeV, and the bosonic charge carriers which render the strings superconducting are squark or slepton condensates. Hence, baryonic or leptonic charge may be stored in the core of the strings. After reviewing the particle physics model under consideration in Section 2, we study the evolution of the string network (Section 3). As we discover, the number density of strings at late times depends sensitively on whether the string loops become vortons or not. In Section 4, we calculate the baryon asymmetry per string loop decay, and use the result to estimate the overall n_B/s ratio generated by this mechanism.

II. LOW ENERGY SUPERSYMMETRY BREAKING MODELS

As we mentioned in the Introduction, supersymmetry provides solutions to many of the puzzles of the standard model such as the stability of the weak scale under radiative corrections as well as the origin of the weak scale itself. Since experimental observations require supersymmetry to be broken, it is essential to have a knowledge of the nature and the scale of supersymmetry breaking in order to have a complete understanding of the physical implications of these theories. At the moment, we lack such an understanding and therefore it is important to explore the various ways in which supersymmetry breaking can arise and study their consequences.

The most common and popular approach is to implement supersymmetry breaking in some hidden sector where some F -term gets a vacuum expectation value (VEV) and then transmits it to the standard model sector by gravitational interactions. This is the so-called hidden $N = 1$ supergravity scenario [15]. If one arranges the parameters in the hidden sector in such a way that the typical $\langle F \rangle$ -term is of the order of $\langle F \rangle^{1/2} \sim \sqrt{(1 \text{ TeV})M_{\text{Pl}}} \sim 10^{11}$ GeV, where $M_{\text{Pl}} = 2.4 \times 10^{18}$ GeV, the gravitino mass $m_{3/2}$ turns out to be of the order of

the TeV scale. In the $N = 1$ supergravity scenario, however, the flavor-blindness is likely to be spoiled in the Kähler potential and one expects large contributions to the Flavour Changing Neutral Currents (FCNC's) at low energies. The suppression of the electric dipole moment of the neutron (EDMN) is also hard to explain in supergravity since one expects the bilinear and trilinear terms to be of the order of $m_{3/2}$ and explicit CP-violating phases to be $\mathcal{O}(1)$.

An alternative to hidden supergravity is provided by the so-called gauge mediated supersymmetry breaking (GMSB) models [16]. The soft supersymmetry breaking mass terms of sfermions get contributions at two-loop, $\widetilde{m}^2 \sim \left(\frac{\alpha}{4\pi}\right)^2 \Lambda^2$, where $\Lambda \sim 10$ TeV is the typical scale of the messenger sector. The GMSB models have the extra advantage that the FCNC effects are naturally suppressed. This is due to the fact that at the scale Λ the squark masses are all degenerate because of the flavor blindness of the standard model gauge group. Only a slight asymmetry is introduced by renormalization group extrapolation from the scale Λ to low energies. Moreover, the trilinear soft breaking terms A vanish at Λ rendering the CP-violation problem milder.

Recently, an alternative approach has been proposed in which gauginos, higgsinos and the third generation of squarks are sufficiently light to stabilize the electroweak scale, but the two first generations of squarks and sleptons are sufficiently heavy to suppress the FCNC and the EDMN below the experimental bound [17–19]. This class of models is dubbed "more" minimally supersymmetric than the MSSM since they do not require some ad hoc supposition of degeneracy or alignment [19]. In most of the attempts made so far along this line [17,18,20,21] the crucial feature is the existence of an anomaly-free $U(1)$ gauge group, which appears anomalous below some scale M . The effective theory includes a Fayet-Iliopoulos (FI) D -term proportional to $\text{Tr } QM^2$, where the Q 's represent the $U(1)$ charges of the heavy fields which have been integrated out. If the first two generations of squarks and sleptons, contrary to the third generation, carry $U(1)$ charges, the required mass hierarchy is obtained. We stress that supersymmetry models can be regarded as realistic only when they are able to reproduce the quantitative features of the fermion spectrum and

the CKM matrix. In a unified picture it is desirable that the solution to all these puzzles reside in the same sector of the theory. Recent investigations have focused on the possibility that a Frogatt-Nielsen flavor mechanism [22] is implemented in the messenger sector [20,21]. This means that the messenger sector is also to be a Frogatt-Nielsen sector and that the messenger $U(1)$ symmetry by which the two first generation fermions get large masses is the same Frogatt-Nielsen $U(1)$ for quarks and leptons. Therefore, the most general (and maybe natural) scenario is an hybrid one where there are two different sources of contributions to ordinary superpartner masses, gauge mediation from the messenger sector and anomalous $U(1)$ D -terms [20,21,23].

As shown in [12], this class of low energy supersymmetry breaking models naturally predicts (superconducting) cosmic strings. Indeed, the gauge group of the secluded sector, where supersymmetry is broken, may be of the form $\mathcal{G} \otimes U(1)_m$. The group $U(1)_m$ is usually some global $U(1)$ which is gauged and made anomaly-free. The fields in the messenger sector are charged under $U(1)_m$ and (some of them) under the standard model gauge group. When supersymmetry is broken in the secluded sector, some scalar fields in this sector may acquire a VEV, but leave the $U(1)_m$ gauge symmetry unbroken. The presence of an anomalous FI D -term ξ may induce the spontaneous breakdown of the residual $U(1)_m$ gauge symmetry along some field direction in the messenger sector. In this case local cosmic strings are formed [12]. Indeed, integrating out the heavy fields belonging to the secluded sector amounts to generate a (say) positive one-loop FI D -term ξ and negative two-loop soft supersymmetry breaking squared masses for (some of) the scalar fields ϕ_i of the messenger sector. These squared masses are generally proportional to the square of the corresponding $U(1)_m$ charges q_i . The potential will read

$$V = - \sum_i q_i^2 \widetilde{m}^2 |\phi_i|^2 + \sum_i \left| \frac{\partial W}{\partial \phi_i} \right|^2 + \frac{g^2}{2} \left(\sum_i q_i |\phi_i|^2 + \xi \right)^2, \quad (1)$$

where g is the $U(1)_m$ gauge coupling constant. From the minimization of (1), one can see that among the fields with F -flat directions, the one with the smallest negative charge (call it φ) will get a VEV

$$\langle |\varphi|^2 \rangle = \frac{1}{q_\varphi} \left(q_\varphi \frac{\widetilde{m}^2}{g^2} - \xi \right). \quad (2)$$

When this happens, the residual global $U(1)_m$ symmetry gets broken leading to the formation of local strings whose mass per unit length is given by $\mu \sim \xi$. Since $\sqrt{\xi}$ is a few orders of magnitude larger than the weak scale, cosmic strings are not very heavy.

We now suppose that the quark and/or the lepton superfields are charged under the $U(1)_m$ group. As we argued, this is in the spirit of the "more" MSSM [19] and somehow welcome when trying to quantitatively predict the fermion mass spectrum and the CKM matrix. Let us focus of one of the sfermion fields, \widetilde{f} with $U(1)_m$ -charge q_f such that $\text{sign } q_f = \text{sign } q_\varphi$. The potential for the fields φ and \widetilde{f} is written as

$$V(\widetilde{f}, \varphi) = -q_\varphi^2 \widetilde{m}^2 |\varphi|^2 - q_f^2 \widetilde{m}^2 |\widetilde{f}|^2 + \frac{g^2}{2} \left(q_\varphi |\varphi|^2 + q_f |\widetilde{f}|^2 + \xi \right)^2 + \lambda |\widetilde{f}|^4, \quad (3)$$

where we have assumed, for simplicity, that \widetilde{f} is F -flat. The parameter λ is generated from the standard model gauge group D -terms and vanishes if we take \widetilde{f} to denote a family of fields parametrizing a D -flat direction.

At the global minimum $\langle \widetilde{f} \rangle = 0$ and the electric charge, the baryon and/or the lepton numbers are conserved. The soft breaking mass term for the sfermion reads

$$\Delta m_{\widetilde{f}}^2 = q_f (q_\varphi - q_f) \widetilde{m}^2, \quad (4)$$

which is positive in virtue of the hierarchy $q_\varphi < q_f < 0$. Consistency with experimental bound requires $\Delta m_{\widetilde{f}}^2$ to be of the order of $(20 \text{ TeV})^2$ or so, which in turn requires $\xi \sim (4\pi/g^2)\widetilde{m}^2 \sim (10^2 \text{ TeV})^2$. Notice that $\Delta m_{\widetilde{f}}^2$ does not depend upon ξ .

Let us analyse what happens in the core of the string. In this region of space, the vacuum expectation value of the field vanishes, $\langle |\varphi| \rangle = 0$, and nonzero values of $\langle |\widetilde{f}| \rangle$ are energetically preferred in the string core

$$\langle |\widetilde{f}|^2 \rangle = \frac{\widetilde{m}^2 q_f^2 - g^2 \xi q_f}{g^2 q_f^2 + 2\lambda}. \quad (5)$$

Since the vortex is cylindrically symmetric around the z -axis, the condensate will be of the form $\widetilde{f} = \widetilde{f}_0(r, \theta) e^{i\eta_f(z,t)}$ where r and θ are the polar coordinate in the (x, y) -plane. One can

easily check that the kinetic term for \tilde{f} also allows a nonzero value of \tilde{f} in the string and therefore one expects the existence of bosonic charge carriers inside the strings. The latter are, therefore, superconducting [24].

What is noticeable is that, besides the electric charge, the baryon/and or the lepton numbers are also broken inside the string. In particular, baryonic charge may be stored in the core of the string. Indeed, if the sfermion particles are identified with some squarks \tilde{q} (eventually parametrizing a standard model D -flat direction), they carry a $U(1)$ baryonic global charge which is derived from the conserved current

$$J_B^\mu = \frac{i}{2} \sum_q q_B^q \left(\tilde{q}^\dagger \partial^\mu \tilde{q} - \tilde{q} \partial^\mu \tilde{q}^\dagger \right), \quad (6)$$

where q_B^q is the baryonic charge associated to any field \tilde{q} . Under the assumption of cylindrical symmetry, the baryonic charge per unit length Q_B along the z -axis will be given by $Q_B = \int d\theta dr r j_B(\theta, r)$ where j_B is the current per unit length along the same axis.

III. STRING NETWORK EVOLUTION

In the class of models considered in Section 2, a network of strings forms during supersymmetry breaking at a temperature $T_c \sim 10^2$ TeV. String formation has been studied extensively [25–29] both analytically and numerically. According to these studies, the initial string correlation length (mean separation between strings) is

$$\xi(t_c) \sim \lambda^{-1} \eta^{-1}, \quad (7)$$

$\eta \sim \sqrt{\xi}$ being the energy scale of the string, and $\lambda \sim g^2$ being a typical coupling constant of the $U(1)$ order parameter field.

The evolution of a string network can be divided into two periods. In the initial period following the string-producing phase transition, the string dynamics is friction-dominated [30–32]. During this period, the correlation length $\xi(t)$ increases super-luminally and catches up to the Hubble radius. The friction-dominated phase ends once $\xi(t) \sim t$. This happens at a time t_f given by

$$t_f \sim (G\mu)^{-1}t_c, \quad (8)$$

G being Newton's gravitational constant and $\mu = \eta^2$ the string mass per unit length. After t_f , the string network enters the scaling epoch during which $\xi(t) \sim t$. Since in our class of models $G\mu \sim 10^{-28}$, all of the physics relevant for baryogenesis takes place deep in the friction-dominated period.

It is important to know the rate at which $\xi(t)$ increases during the friction-dominated period. We take this rate to be given by

$$\xi(t) = \xi(t_c) \left(\frac{t}{t_c} \right)^{3/2}. \quad (9)$$

This result can be justified as follows. For a string loop, tension dominates over friction if the loop radius R is larger than a critical radius $r_f(t)$ which increases in time since the friction per unit string length decreases. A comparison between tension and friction gives [30–32]

$$r_f(t) = G\mu M_{\text{Pl}}^{1/2} t^{3/2}. \quad (10)$$

If $\xi(t)$ were smaller than $r_f(t)$, then the string network would evolve under the force of tension alone, and the resulting intercommutations of string would lead to a rapid increase in $\xi(t)$ [33]. Conversely, if $\xi(t) \gg r_f(t)$, then strings would be approximately static in comoving coordinates, which would lead to a decrease of the ratio $\xi(t)/r_f(t)$. Hence,

$$\xi(t) \sim r_f(t) \quad (11)$$

is a dynamical fixed point of the string network in the early period of evolution.

The rapid increase of $\xi(t)$ given by (9) is achieved by string loop production. Unless small-scale structure on the string network is important already in the friction-dominated epoch, the loops are produced with a radius proportional to $\xi(t)$ (see the reviews in [34–36]). In this case, the number density of loops created per unit time is

$$\frac{dn}{dt} = \nu \xi^{-4} \frac{d\xi}{dt}, \quad (12)$$

where ν is a constant of order unity.

In the class of models considered, the strings are superconducting and carry baryon charge. Let Q_B denote the charge per unit length. By order of magnitude considerations, the value of Q_B will be of the order η . Then, the charge on a correlation length of string at the time t_c is

$$Q_c = Q\xi(t_c). \quad (13)$$

On longer pieces of string, the charge adds up as a random walk [37–39]. Hence, loops created at time t with radius $\xi(t)$ have a root mean square charge of

$$Q_B(t) = \left[\frac{\xi(t)}{\xi(t_c)} \right]^{1/2} Q_c. \quad (14)$$

Whereas non-superconducting string loops decay only gravitationally with power $P_G = \gamma G\mu^2$, γ being a constant whose value is about 100 [35], strings with a current on them decay predominantly by electromagnetic radiation. The power P_γ of electromagnetic emission can be parametrized as [40]

$$P_\gamma = \gamma_{em} j^2 \mu. \quad (15)$$

Here, γ_{em} is a constant analogous to γ (and for the numerical estimates which follow we will set them equal), and the relative current j is defined by

$$j = \frac{J}{J_{max}}, \quad J_{max} = eg_e^{-1}\eta, \quad (16)$$

where J denotes the current on the loop, and g_e is the electromagnetic gauge coupling constant. Note that the current increases as the radius R of the loop decreases from its initial value R_0 :

$$j(R) = j_0 \left(\frac{R_0}{R} \right), \quad (17)$$

j_0 being the initial relative current. Comparing the power of gravitational and electromagnetic radiation, it follows that if

$$j_0 > (G\mu)^{1/2}, \quad (18)$$

then electromagnetic radiation is the dominant energy loss mechanism during the entire life-time of the loop. Note that

$$\dot{R} = (\beta\mu)^{-1}P, \quad (19)$$

where P is the total power of radiation and βR is the length of the loop (in general, $\beta \neq 2\pi$ if the loop is not a circle). Combining (15), (17) and (19), we can solve for the life-time of a string loop. We find that provided

$$j_0 > (G\mu)^{1/4}(R_0\eta)^{1/2}, \quad (20)$$

then the life-time of the loop is less than one Hubble expansion time.

The current per correlation length on the string is of the order η (by dimensional arguments). On longer string segments, the current adds up as a random walk (analogous to (14)). Hence,

$$j_0(R) \sim \left[\frac{R}{\xi(t_c)} \right]^{1/2} eg_e^{-1}, \quad (21)$$

and hence the conditions (18) and (20) are clearly satisfied. Thus, on cosmological time-scales, superconducting strings instantaneously collapse. This fact will simplify the computation of the baryon asymmetry.

As mentioned in the Introduction, the cosmic string network provides a means for trapping baryon charge until late times. In the following section we will discuss how the charge trapped on a string loop gets converted into a net baryon number when the loop decays. For now we simply write

$$\Delta n_B = Q_B \epsilon, \quad (22)$$

for the net baryon number generated by the decay of a loop of charge Q_B , where ϵ is a constant related to the CP-violation parameter (see Section 4).

The resulting baryon asymmetry depends crucially on whether the loops decay immediately or collapse to form vortons [41], which themselves decay after t_{EW} . In the former case, the only loops which contribute to the net n_B/s are those produced after t_{EW} :

$$n_B(t) = \int_{t_{EW}}^t dt' \epsilon Q_B(t') \frac{dn}{dt}(t') \left(\frac{t'}{t}\right)^{3/2}, \quad (23)$$

where the final factor is due to the cosmological redshift. The integral (23) is dominated by the contribution from near t_{EW} (or, more generally, the time when sphaleron processes fall out of equilibrium). Inserting (9), (12) and (14), the integral becomes

$$n_B(t) \simeq \frac{2}{3} \epsilon \nu Q_c \xi(t_c)^{-3} \left(\frac{t_c}{t}\right)^{3/2} \left(\frac{t_c}{t_{EW}}\right)^{9/4}. \quad (24)$$

The last factor in (24) is a geometrical suppression factor resulting since baryons produced by loops which decay before t_{EW} equilibrate by sphaleron processes.

On the other hand, if superconducting loops form vortons which decay only after t_{EW} , then all loops created since t_c contribute to the net n_B which is determined now by an integral like (23) but with lower integration limit t_c . Since the integral is once again dominated by the lower limit we obtain

$$n_B(t) \simeq \epsilon \nu Q_c \xi(t_c)^{-3} \left(\frac{t_c}{t}\right)^{3/2}. \quad (25)$$

The geometric suppression factor which appeared in (24) is no longer present. For $T_c \sim 10^2$ TeV, baryogenesis is more efficient by a factor of $(T_c/T_{EW})^{4.5} \sim 10^9$ if the loops form vortons which are stable until at least t_{EW} . Obviously, the overall strength of our baryogenesis mechanism depends on the value of ϵ , to the evaluation of which we now turn.

IV. THE BARYOGENESIS MECHANISM

As mentioned in the Introduction and discussed in Section 2, the strings in our class of theories carry baryonic charge in the form of a squark or slepton condensate. When the string loops decay, a net baryon charge is generated analogously to the Affleck-Dine mechanism for baryon number production [42]. We should comment here on two crucial features of

the mechanism. First, when superconducting strings are formed, one may expect that the number of strings with some baryonic charge Q_B is equal to the number of strings with opposite baryonic charge $-Q_B$. Therefore, when string loops decay and release the baryon number in the thermal bath, the net baryon asymmetry vanishes unless some source of *explicit* CP-violation is present in the interactions. Secondly, even though superconducting strings are formed at the temperature $T_c \gg T_{EW}$, the baryon charge stored in the core is not washed out by baryon number violating processes induced by sphaleron transitions. Indeed, sphalerons are not active in the strings. This is because *any* condensate ϕ_i carrying $SU(2)_L$ quantum numbers contributes to the sphaleron energy, $E_{\text{sph}} \propto \sqrt{\sum_i |\phi_i|^2}$. Therefore, if the sfermion condensate in the core of the string is charged under $SU(2)_L$, sphaleron transitions are suppressed inside the strings (see also [43]). This is equivalent to saying that the superconducting strings act like “bags” containing baryon charge and protect it from sphaleron wash-out throughout the evolution of the Universe until baryon number violating processes are rendered harmless.

For sake of concreteness, let consider the case in which the sfermion condensate is formed by the third family left-handed stop \tilde{t}_L (one can easily generalize this case). The left-handed stop decay channels most important for our considerations are $\tilde{t}_L \rightarrow t_R + \widetilde{H}_2^0$, $H_2^0 + \tilde{t}_R$ where H_2^0 , \widetilde{H}_2^0 and \tilde{t}_R are the Higgs, the Higgsino and the right-handed stop, respectively. The constant ϵ (introduced in (22)) which determines the asymmetry in baryon and antibaryon production is determined by the difference in the decay widths between \tilde{t}_L and \tilde{t}_L^* :

$$\epsilon = \frac{\Gamma_{\tilde{t}_L} - \Gamma_{\tilde{t}_L^*}}{\Gamma_{\tilde{t}_L} + \Gamma_{\tilde{t}_L^*}}. \quad (26)$$

The value of ϵ is determined by the strength of the explicit CP-violation which couples to the \tilde{t}_L decay processes. For us, the dominant contribution is the CP-violating phase of the coefficient A_t of the term

$$A_t h_t \tilde{t}_L \tilde{t}_R^* H_0^2 + \text{h.c.} \quad (27)$$

in the interaction Lagrangian, where h_t is top Yukawa coupling.

The contribution to ϵ comes from the interference terms between the tree level and one loop decay diagrams. Assuming that $m_t \simeq \widetilde{m}$ (the supersymmetry breaking mass) we obtain

$$\epsilon \simeq \frac{\ln 2}{8\pi} h_t^2 \sin\phi_{A_t}, \quad (28)$$

where we have written $A_t = |A_t| e^{i\phi_{A_t}}$. Notice that ϕ_{A_t} is not constrained from the experimental upper bound on the EDMN and therefore can be as large as unity.

At this point, we can insert the value of ϵ obtained above into the general expressions for the baryon number obtained in Section 3. Replacing time t by the corresponding temperature T by means of the Friedmann equations, and making use of the expression for the entropy density s in terms of the number of degrees of freedom g_* in thermal equilibrium at T_{EW}

$$s(T) \sim g_* T^3, \quad (29)$$

we obtain from (24) and (7):

$$\frac{n_B}{s} \sim \epsilon \nu Q_c \lambda^3 (g_*)^{-1} \left(\frac{T_{EW}}{T_c} \right)^{4.5} \quad (30)$$

for the case when the string loops collapse and disappear immediately, and from (25) and (7)

$$\frac{n_B}{s} \sim \epsilon \nu Q_c \lambda^3 (g_*)^{-1} \quad (31)$$

for the case when stable vortons form during the initial loop collapse and survive until after t_{EW} .

We conclude that in the case in which stable vortons form upon the collapse of a string loop, it is possible to obtain the observed value of n_B/s without unnatural constraints on the coupling constants. To get the feeling with the numbers, for $\epsilon \sim 10^{-2}$, $\lambda \sim 10^{-1}$, $\nu \sim 1$ and $g_* \sim 10^2$, we need $Q_c \sim 10^{-3}$ to produce $n_B/s \sim 10^{-10}$.

V. CONCLUSIONS

In this paper we have proposed a new mechanism for the generation of the baryon asymmetry in the early Universe. It is based on the general observation that, if superconducting cosmic strings form at scales within a few orders of magnitude of the weak scale and carry some baryon charge, the latter is efficiently preserved from the sphaleron erasure and may be released in the thermal bath at low temperatures. A natural framework for this mechanism is represented by the class of models where supersymmetry is broken at low energy. In such a case, the charge carriers inside the strings are provided by the scalar superpartner of the fermions which carry baryon (lepton) number. Since these scalar condensates are charged under $SU(2)_L$, baryon number violating processes are frozen in the core of the strings and the baryon charge number can not be wiped out at temperatures larger than T_{EW} . The mechanism is very efficient if stable vortons form upon the collapse of the string loops and survive until after the electroweak phase transition (note that the role of vortons in baryogenesis has recently also been studied in [44]).

We like to stress that our proposal has a number of advantages with respect to the idea of bubble-mediated electroweak baryogenesis since it does not ask for a sufficiently strong electroweak first order phase transition and therefore does not impose any stringent bound on the mass spectrum of the MSSM. An important issue for future study is to investigate the probability of formation of vortons and their stability. As has been shown recently [45], cosmological constraints from the density parameter at the present time and from nucleosynthesis are consistent with vorton formation at the scale considered in this paper.

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