BARYON NUMBER PENETRABILITY AS A MEASURE OF ISOTHERMAL BARYON NUMBER FLUCTUATIONS IN THE EARLY UNIVERSE

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Abstract

We have examined the efficiency of baryon-number transport mechanism across the phase boundary in a cosmological quark-hadron phase transition through the proper estimate of baryon-number penetrability $\langle \Sigma_h \rangle$. For this purpose we have derived first the double-pair creation probability $P_b$ in terms of single-pair creation probability per unit time and unit volume $\kappa_m$ and then obtained an analytical expression for $\langle \Sigma_h \rangle$. Our calculation is free from the uncertainty of the value of double-pair creation probability per unit time and unit volume $\kappa_b$ which was used as a free parameter in earlier calculations. Finally the variations of double-pair creation probability $P_b$ as well as $\langle \Sigma_h \rangle$ with temperature are shown and compared with other known results.

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1 Introduction

Recently there is a growing interest in the primordial nucleosynthesis in an Universe having an inhomogeneous distribution of baryon number. The most popular mechanism for the generation of baryon number inhomogeneities has been suggested as a first-order QCD phase transition [1-8]. The present Universe is presumed to have undergone several successive phase transitions associated with symmetry breaking in its early stage. A QCD phase transition might have occurred at approximately $10^{-6}$ sec after the big bang in the early Universe. Even though many aspects of phase transition are not a settled issue, yet, the QCD phase transition is expected to be a first-order one from lattice QCD results [9]. Therefore, a large local fluctuation in the baryon to photon ratio $n_B/n_\gamma$ could arise which might subsequently modify the standard picture of the primordial nucleosynthesis (PNS). Such a fluctuation might also result into the formation of quark nuggets and/or quark stars etc.[1-8]. The basis for the production of isothermal baryon number fluctuation lies in the separation of cosmic phases as follows [1,5]. Initially the Universe is in the quark-gluon plasma (QGP) phase at a high temperature and the net baryon number resides entirely in the quark-gluon plasma phase and is distributed homogeneously. As the Universe expands, the temperature drops to the critical temperature $T_c$ where the quark-gluon plasma exists in thermal and chemical equilibrium with the dense and hot hadron gas. Subsequently, the expansion requires a continuous conversion of QGP into the hadron phase. The phase transition is completed when all the quark-gluon plasma has been converted to hadron phase and all the baryon number residing finally in the hadron phase is distributed homogeneously. The magnitude of baryon number fluctuation is estimated by the baryon contrast ratio $R_{eq}$ of the net equilibrium baryon number density in the quark-gluon plasma (QGP) phase to that in the hadron gas (HG) phase, i.e., $R_{eq} = n_B^{QGP}/n_B^{HG}$ which is evaluated at the critical temperature and chemical potential. The baryon number density inhomogeneity arising due to such a quark-hadron phase transition will thus alter the yields of the PNS which is a function of $R_{eq}$ and the neutron to proton ($n/p$) ratio [10]. Several theoretical
attempts have been recently made to determine the values of $R_{eq}$ [11-14].

All aforementioned calculations of $R_{eq}$ assume that the Universe is in complete thermal and chemical equilibrium during the phase transition [15]. But in reality, the cosmic first-order phase transition will necessarily result in deviations from thermal and chemical equilibrium because the low temperature hadron phase is not nucleated immediately at the critical temperature $T_c$. A generic feature of quantum or thermal nucleation theory is that the nucleation rate does not become large unless the temperature has dropped below $T_c$ [16]. The magnitude of these deviations will depend on the efficiency of heat and baryon transport across the transition front during the phase transition. Latent heat transport is carried out by the motion of the boundary wall which converts volume of one vacuum into another. This vacuum energy difference acts as a source for particle creation. However, the latent heat or entropy could also be carried out across the phase boundary by neutrinos apart from surface evaporation of hadrons (mostly pions). Baryon number is not thermally created in the hadron phase. Therefore, it actually flows across the boundary by the strong interaction physics if the hadron phase is going to have any net baryon number at all.

There are two limiting situations governing the efficiency of baryon number transport across the phase boundary [5.17]. In the first situation, when the bubble-like regions of quark-gluon plasma are small in size, the phase boundaries move slowly compared to the baryon number transport rate. Hence the baryon number is quickly and efficiently transported across the boundaries, so as to maintain the chemical equilibrium. In the second situation, the boundaries of the bubble-like regions move rapidly compared to the baryon transport rate so that the net baryon transport process becomes inefficient. Therefore, chemical equilibrium may not be achieved on the time scale of the phase coexistence evolution and, consequently, the baryons could be concentrated in the shrinking bubble-like regions of the QGP. The magnitude of the resulting baryon number fluctuation at the time of phase decoupling will still be given by the ratio $R'$ (say) of baryon number densities in the two phases but now each phase will have a different baryon
chemical potentials. Thus, the final ratio $R'$ will be larger than the equilibrium ratio $R_{eq}$.

It is clear that in both the abovementioned situations isothermal baryon number fluctuations will result. However, the magnitude of $R_{eq}$ is in the limit of efficient baryon number transport mechanism which is the aforesaid first situation. The efficiency of baryon number transport across the phase interface will be given by the bulk properties of the phases and the estimate of baryon number penetrability [5,17]. The baryon number penetrability $\Sigma_h$ from the hadronic side is defined as the probability that a baryon which approaches from the hadronic phase will dissociate into quarks and pass over into the QGP phase. Similarly from another side, we could define the probability $\Sigma_q$ that the quarks which approaches towards the phase boundary will form a color singlet baryon and pass over to the hadron phase. The thermal average of these probabilities are related by the detailed balance

$$k \ f_{q-\bar{q}}(\Sigma_q) = f_{b-\bar{b}}(\Sigma_h)$$

where $f_{q-\bar{q}}$ ($f_{b-\bar{b}}$) is the excess quark flux (baryon flux) in the QGP (HG) phase, respectively, and the dimensionless quantity $k$ takes values from $1/3$ to $1$. The value $k = 1/3$ implies that baryon number predominantly passes over the phase boundary as preformed three-quark clusters in the QGP phase. Similarly $k = 1$ signifies that baryons will be predominantly formed by the leading quarks which cross over the phase boundary into the hadron phase. A low value for $\langle \Sigma_h \rangle$ (or $\langle \Sigma_q \rangle$) ($<1$) implies an early drop out of chemical equilibrium during the separation of phases and thus might lead to the large amplitude of baryon-number fluctuations.

The estimation of baryon number penetrability has been attempted recently [5,17,18] by many authors within the framework of the chromoelectric flux tube (CFT) model. Here an energetic leading quark in the QGP is assumed to pick up an antiquark forming thereby an expanding CFT whose decay through a single or double pair creation results in the formation of mesons or baryons, respectively. Naturally, the calculation of $\langle \Sigma_h \rangle$ within CFT model requires the values of the single as well as coherent double pair cre-
ation probabilities per unit space per unit time, called $\kappa_m$ and $\kappa_b$, respectively. Although $\kappa_b$ is a crucial factor for the formation of a baryon yet no theoretical expression for it based on strong interaction physics exists in the literature.

In this context the following points pertaining to a recent paper by Jedamzik and Fuller [17, referred to as JF hereafter] become particularly relevant: (i) JF have extracted the numerical value of $\kappa_b$ by analyzing empirically the ratio of baryons to mesons produced in the $e^+e^- \rightarrow$ hadrons experiment conducted in the accelerator laboratories. (ii) JF allow for the full angular range $0 \leq \theta \leq \pi$ where $\theta$ is the polar angle of the leading quark’s velocity with respect to the normal to the phase boundary. (iii) The final expression for $\langle \Sigma_h \rangle$ written as an integral over the quark distribution functions is computed by JF entirely numerically. (iv) Finally, the threshold energy in JF’s work increases with the temperature in an unreasonable manner.

The aim of the present paper is to extend/modify the theory of JF [17] in multifold respects: (i) By regarding the unconnected double-pair creation event as a succession of two single-pair events, we show in Sec.2 below, that the double-pair creation probability $P_b$ (over finite time duration) can be obtained in terms of the $\kappa_m$ parameter itself. (ii) Since the leading quark’s velocity should be directed towards the phase boundary, we point out in Sec.3 that the corresponding polar angle should be restricted to the range $0 \leq \theta \leq \pi/2$. (iii) We derive in Sec.3 an analytical expression for the thermal average of the baryon number penetrability so that the dependence of $\langle \Sigma_h \rangle$ on the relevant variables becomes more transparent. (iv) We show in Sec.4 the comparison between our and JF models by plotting $P_b$ as well as $\langle \Sigma_h \rangle$ as functions of the temperature when the quarks are assigned either the constituent or current mass. (v) Finally, we carefully analyze the temperature-dependence of the threshold energy in Sec.4 to obtain an elegant formula for $\langle \Sigma_h \rangle$ in the high T limit.
2 Decay Statistics of Flux Tubes

Chromoelectric flux tube (CFT) models provide us a phenomenological description regarding the formation of a hadron from quarks and gluons [19]. These models assume the existence of a chromoelectric field between two oppositely coloured quarks where the chromoelectric field strength is assumed to be constant in magnitude and independent of the separation of the quarks. These fields can be thought of as being confined to a tube of constant width known as flux tube. Lattice QCD results justify the existence of such a flux tube [21]. Initially, chromoelectric flux tube models were used to explain the processes such as $e^+e^- \rightarrow \text{hadrons}$ [22,23]. Later these models have also been used to describe the spectrum of mesonic and baryonic resonances [22,24]. Recently flux tube models have been employed to estimate the meson evaporation rates from quark-gluon plasma produced in relativistic heavy-ion collisions [20] and also to estimate baryon-number penetrability across the phase boundary in cosmological quark-hadron phase transition [17,18].

If the CFT is regarded as an unstable system then its stochastic properties can be conveniently discussed by introducing the following abbreviations : SP $\equiv$ survival probability, DP $\equiv$ decay probability, m $\equiv$ mesonic or single-pair production channel, b $\equiv$ baryonic or connected double-pair production channel, $b^*$ $\equiv$ disconnected double-pair production channel. Let us now take up the the earlier model used by Jedamzik and Fuller [17] and a new proposal due to us.

2.1 Jedamzik and Fuller (JF) Model

In analogy with QED, at a perturbative field theoretic level, the JF model amounts to connected diagrams shown in Figs. 1(m) and 1(b). The corresponding DP’s per unit time per unit volume are denoted by $\kappa_m$ and $\kappa_b$, respectively, having a total $\kappa = \kappa_m + \kappa_b$. 
Figure 1: (m), (b), (b*): Diagrams drawn in analogy with perturbative QED, representing decay into single, connected double, and unconnected double $q\bar{q}$ pair channels, respectively. The crosses stand for the chromoelectric field as a source. The dotted lines denote gluons.

The space-time integrated parameters are called as

\[
\begin{align*}
    w_m &= \kappa_m \int_0^{t_0} V(t)dt \\
    w_b &= \kappa_b \int_0^{t_0} V(t)dt
\end{align*}
\]

and $w = w_m + w_b$, $V(t)$ is the instantaneous volume of the flux tube and $t_0 \sim E/\sigma$ is the maximum time upto which it expands when the incident quark has energy $E$. From the conservation of energy and momentum parallel to the phase boundary, the above integral could be written as in [17]

\[
\int_0^{t_0} V(t)dt = \frac{\pi \Lambda^2}{2\sigma^2} E^2 \cos \theta
\]

with $\theta$ being the angle of incidence of the leading quark relative to the normal to the phase boundary, $\Lambda$, the radius of the flux tube and $\sigma$ the string tension of the flux tube.
Then, according to the classical theory of radioactivity, the net SP and DP counting all channels would be

\[ P_s = e^{-w} \]  
\[ P_d = 1 - e^{-w} \]  

Since the relative weights for the \(m\) and \(b\) channels are \(w_m/w\) and \(w_b/w\), respectively, the corresponding channelwise DP’s are as used by JF:

\[ P_m = \frac{\kappa_m}{\kappa} P_d \]
\[ = \frac{\kappa_m}{\kappa} \left( 1 - \exp \left[ -\kappa \int_0^{t_0} V(t) dt \right] \right) \]  
\[ P_b = \frac{\kappa_b}{\kappa} P_d \]
\[ = \frac{\kappa_b}{\kappa} \left( 1 - \exp \left[ -\kappa \int_0^{t_0} V(t) dt \right] \right) \]  

Although \(\kappa_m\) can be estimated from Schwinger’s single-pair production mechanism [25] applied to QCD, nothing a priori is known about \(\kappa_b\). However, JF have extracted empirically the magnitude of \(\kappa_b\) from the analysis of \(e^+e^- \rightarrow \text{hadrons}\) data.

### 2.2 Our Proposal

Suppose we retain information about the single \(q\bar{q}\) pair production of Fig. 1(m) but ignore the diagram 1(b). Then, the double-pair production event may be looked upon as a sequence of two unconnected single-pair creations within time \(t_0\) as shown in Fig.1(b*).

Since the flux tube is now characterized only by the parameter \(w_m\), hence radioactivity theory would give the following expressions for the net SP and DP over the duration time \(t_0\):

\[ P_s^* = e^{-w_m} \]  
\[ P_d^* = 1 - e^{-w_m} = \sum_{n=1}^{\infty} p_n^* \]  

where the DP of having exactly \(n\) successive single-pair events is a Poissonian viz.

\[ p_n^* = e^{-w_m} \frac{w_m^n}{n!} \]
Therefore, in our approach the DP for the single and double-pair would be obtained from

\[ P_m^* = p_1^* = e^{-w_m} w_m \quad (12) \]

\[ P_b^* = p_2^* = e^{-w_m} \frac{w_m^2}{2} \quad (13) \]

In contrast to the JF approach \((7, 8)\), our proposal involves only one parameter \(w_m\) (or \(\kappa_m\)) as an input.

### 3 Analytical Estimation of Baryon number Penetrability

The main result of JF’s work is contained in their baryon penetrability equation which gives the thermal average of the baryon penetrability integrated over the \(q\bar{q}\) distribution function. We wish to make two plausible comments on this result. Firstly, the angular integration range in their calculation \(0 \leq \theta \leq \pi\) may not be appropriate because if the leading quark is to enter the phase boundary then it should have \(0 \leq \theta \leq \pi/2\). Therefore we suggest that the corrected expression should be

\[ \langle \Sigma_h \rangle = \frac{1}{f_{b-\bar{b}}} \int_0^{\pi/2} d\theta \int_{E_{th}}^{\infty} dE \frac{dn_{q-\bar{q}}}{dEd\theta} \dot{x}_q^q(\theta) P_b(E, \theta) \quad (14) \]

Here \(f_{b-\bar{b}}\), the excess baryon flux in the hadron phase, is given by \([17]\)

\[ f_{b-\bar{b}} = \frac{\mu_b}{T} \frac{g_b}{2\pi^2} (m_b + T) T^2 \exp(-m_b/T) \quad (15) \]

where \(m_b, \mu_b\) and \(g_b\) are the baryon mass, baryon chemical potential and degeneracy factors, respectively and \(\dot{x}_q^q\) is the component of the leading quark velocity perpendicular to the phase boundary, and \(E_{th}\) is the threshold energy.

Secondly, JF have done their subsequent calculations numerically so that the dependence of \(\langle \Sigma_h \rangle\) on the relevant parameters remains somewhat obscure. We suggest that an analytical estimate for the seemingly complicated expression \((14)\) is very desirable because
that would make the dependence of $\langle \Sigma_h \rangle$ on the relevant parameters more transparent.

For this purpose we proceed as follows:

Using Boltzmann approximation ($E/T \gg 1$), the differential excess quark number density in a given energy interval $dE$ and in a given interval of incident angles $d\theta$ becomes

$$\frac{dn_{q-\pi}}{dEd\theta} \sim \frac{\mu_q}{T} \frac{g_q}{2\pi^2} E^2 \exp(-E/T) \sin \theta$$

(16)

where $\mu_q, g_q$ are the quark chemical potential, and the statistical weight of quarks, respectively.

Then Eq.(14) becomes

$$\langle \Sigma_h \rangle \approx C \int_0^1 ds \int_{E_{th}}^\infty dE \ E^2 \exp(-E/T) \ P_b(E, \theta)$$

(17)

where

$$C \equiv \frac{1}{f_{b-\bar{b}}} \frac{\mu_q}{T} \frac{g_q}{2\pi^2}; \ s = \cos \theta$$

(18)

Due to the rapidly damped factor $\exp(-E/T)$ in Eq.(17) most contribution to the energy integration comes from around the threshold energy $E_{th}$. Therefore, it is convenient to make the transformation

$$\rho = \frac{E - E_{th}}{T}; \ E = E_{th}\{1 + \frac{T}{E_{th}}\rho\}$$

(19)

Upon using the identity $\int_0^\infty d\rho \ exp(-\rho) = 1$, Eq.(17) reduces to

$$\langle \Sigma_h \rangle \approx C \int_0^1 ds \ T \ E_{th}^2 \exp(-\frac{E_{th}}{T}) \ P_b(E_{th}, \theta)$$

(20)

within correction terms of order $T/E_{th}$. Now, at large angles $\theta \to \pi/2$, $s \to 0$, $E_{th} \to \infty$ implying that $\exp(-E_{th}/T)$ gets heavily damped. Hence, most contribution to the angular integral must come from around the forward direction $\theta \to 0$, $s \to 1$. For convenience define

$$\hat{E}_{th} \equiv E_{th}|_{\theta=0}$$

$$= m_b + Bn_q^{-1}$$

(21)
\[ P_b \equiv P_b(E_{th}, \theta = 0) \]  
\[ \lambda \equiv \frac{E_{th} - \hat{E}_{th}}{T} \approx \frac{m_b}{T} (1 - s) \]  

Here \( B \) is the bag constant and \( n_q \) is the quark plus antiquark density. The above eq.(21) is the threshold energy condition in the forward angle direction (\( \theta = 0 \)) and the next eq.(22) is the corresponding probability. Using the identity \( \int_0^{m_b/T} d\lambda \exp(-\lambda) \approx 1 \), eq.(20) yields

\[ \langle \Sigma_h \rangle \approx C T^2 \frac{E_{th}^2}{m_b} \exp(-\hat{E}_{th}/T) \hat{P}_b \{ 1 + O(T/m_b) \} \]  

Substituting Eq.(18) for \( C \), we arrive at the desired analytical estimate :

\[ \langle \Sigma_h \rangle = \frac{\mu_q}{\mu_b} \frac{g_q}{g_b} \left( \frac{E_{th}}{m_b} \right)^2 e^{-\hat{E}_{th}/T} \hat{P}_b \{ 1 + O(T/m_b) \} \]  

The probability function in JF model

\[ \hat{P}_b = a \left[ 1 - \exp\{-b \hat{E}_{th}^2\} \right] \]  

can be replaced in our model as

\[ \hat{P}_b^* = e^{-w_m} \frac{w_m^2}{2} \]  

with

\[ a = \frac{\kappa_b}{\kappa_m + \kappa_b} ; \quad b = (\kappa_m + \kappa_b) \frac{\pi \Lambda^2}{2\sigma^2} \]  
\[ w_m = \kappa_m \frac{\pi \Lambda^2 \hat{E}_{th}^2}{2\sigma^2} \]  

Here the string tension of the flux tube \( \sigma \simeq 0.177 GeV^2 \), and single-pair creation probability per unit time per unit volume \( \kappa_m \) is given by [25]

\[ \kappa_m = \frac{\sigma^2}{4\pi^3} \exp(-\pi m_q^2 / \sigma) \]  

where \( m_q \) is the quark mass. Since there is an ambiguity in choosing the precise value of \( m_q \), we allow in the sequel both the possibilities viz. \( m_q = 300 MeV \) (constituent quark mass) and \( m_q = 10 MeV \) (current quark mass). Eqs.(25), (26), (27) form the main algebraic results of the present paper.
4 Results and Discussions

In Fig. 2, we have shown the variation of double pair creation probability at threshold in the forward direction with the temperature for a small value of $b = 0.047 \text{ GeV}^{-2}$ corresponding to the constituent quark mass ($m_q = 330 \text{ MeV}$). Here we have taken the value of the ratio $a(\approx \kappa_m/\kappa_b) \approx 1/5$ as in [17].

Figure 2: Variation of double pair creation probability $\hat{P}_b (\hat{P}_b^*)$ at threshold in the forward direction with the temperature $T$ when the quark has constituent mass. The dashed curve represents our model calculation (cf. Eqs.(13), (27)) whereas the solid line represents the JF model (cf. Eqs.(8), (26)). For remaining symbols see Eqs.(28,29,30).

Fig.3 is also the same as Fig.2 except that a higher value of $b = 0.32 \text{ GeV}^{-2}$ is taken corresponding to the current quark mass ($m_q = 10\text{ MeV}$) and the ratio $a \approx 1/20$.

Clearly, the difference between the predicted values of $\hat{P}_b$ in our model and JF model
must be attributed to their different premises. Our approach has the advantage of working with only one decay parameter $\kappa_m$ but ignores the connected diagram 1(b). The JF approach has the advantage of including the graph 1(b) but at the cost of bringing-in the additional unknown decay parameter $\kappa_b$.

Next, we turn to the calculation of the baryon number penetrability $\langle \Sigma_h \rangle$ based on the analytical result (25). The probability $\hat{P}_b$ can have either the JF form (26) or our form (27) with the remaining parameters having been set as for u, d flavours in QGP sector and nucleons in the hadronic sector. Figs. 4 and 5 display the dependence of $\langle \Sigma_h \rangle$ on the temperature when the quark has the constituent and current masses, respectively.

In Fig.4, when the quark has constituent mass, our curve lies below that of the JF model. Our small predicted value $\langle \Sigma_h \rangle \sim O(10^{-3})$, incidentally, justifies the results.
Figure 4: Variation of baryon number penetrability \( \langle \Sigma_h \rangle \) (cf. Eqs.(25), (26), (27)) with the temperature \( T \) when the constituent quark mass is chosen. Other notations are the same as in Fig.2 obtained by Sumiyoshi et al.[18] and Fuller et al.[5]. However, in the opposite limit, i.e., for current quark mass, Fig.5 reveals that our graph is consistently above that of JF. Our predicted higher value of \( \langle \Sigma_h \rangle \simeq O(10^{-2}) \) seems to agree with the original result obtained by Jedamzik and Fuller [17] using numerical quadrature involving the unphysical angular range \( 0 \leq \theta \leq \pi \).

In passing we note that the typical time upto which the CFT expands is \( t_0 \sim \hat{E}_{th}/\sigma \sim 10^{-24}\text{s} \) which seems to be one order of magnitude smaller than typical nuclear time scale of \( 10^{-23}\text{s} \). Can the \( b^* \) channel of Fig.1 be important under such a circumstance? To
answer this question we borrow the $w_m$ values from Figs. 2-5 and look at the ratio

$$\frac{P_b^*}{P_m^*} = \frac{w_m}{2} \sim 2\% \text{ to } 15\%$$

which is sizable. Therefore, our view of regarding the double-pair production event as a succession of two single-pair events can be justified.

Towards the end we wish to comment on a couple of factors on which baryon number penetrability crucially depends within the chromoelectric flux tube approach. One factor is the choice of the quark mass as seen from Figs. 4 and 5. Another factor, in which $\langle \Sigma_h \rangle$ is very sensitive, is the threshold energy $\hat{E}_{th}$ in the forward direction. Following the idea of Jedamzik and Fuller [17], this threshold energy $E_{th}$ consists of two parts: one is simply the rest mass of the baryon $m_h$ and the other is the interaction energy $Bn_q^{-1}$ for each quark and antiquark which resides in the QGP. JF have parametrized the
interaction energy as $\approx 3.7T$ which apparently grows with the temperature. However, in our opinion, the function $Bn_q^{-1}$ should decrease with $T$ like $T^{-3}$ for a fixed value of the bag constant $B$ since the total quark density $n_q$ is an increasing function of $T$. So $E_{th}$ would tend to $m_b$ in Eq.(25) which yields the elegant high-temperature behaviour

$$\langle \Sigma_b \rangle \overset{T \to \infty}{\to} \hat{P}_b$$

(32)

remembering that $g_q = 12$, $g_b = 4$ and $\mu_b = 3\mu_q$.

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