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EQUILIBRIUM FLAVOR DYNAMICS DURING THE COSMIC CONFINEMENT TRANSITION

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B. Kämpfer: Equilibrium flavor dynamics during the cosmic confinement transition. KFKI-1988-59/A

ABSTRACT

The dynamics of the flavor composition of strongly interacting matter during the cosmic confinement transition is followed in a simplified thermodynamical model. Relying on thermal, mechanical and chemical equilibrium we analyse the strangeness fraction of strongly interacting matter. Due to equilibrium with respect to $\Delta S = 0$ and $\Delta S = 1$ weak interactions the relations between different flavors depend strongly on the poorly known lepton excess. In an universe where the lepton (antilepton) excess is in the same order of magnitude as the baryon excess, the strange quark abundancies are suppressed (enhanced). In the hadron phase the strange baryons carry up to a half of the baryon excess.

Б.Кэмпфер: Динамика ароматов во время космического перехода по конфайнменту. KFKI-1988-59/A

АННОТАЦИЯ

Используется простая динамическая модель для описания динамики ароматов в сильной взаимодействующей космической материи во время перехода к конфайнменту. Исследуются последствия термического, механического и химического равновесия для числа странных частиц. Вес ароматов зависит от числа лептонов и условий равновесия в слабых реакциях, странность меняется на $\Delta S=0$ или $\Delta S=1$. Когда число лептонов во вселенной больше (меньше) числа антилептонов, тогда и существует больше (меньше) числа странных кварков. В адронной фазе эти странные адроны носят до половины барионного заряда.

Kämpfer B.: Egyensúlyi ízdinamika a kozmikus kvarkbezáró átmenetben. KFKI-1988-59/A

KIVONAT

Egyszerűsített termodinamikai modellben nyomonkövetjük a kozmikus kvarkbezáró fázisátmenet során az erősen kölcsönható anyag ízösszetételének dinamikáját. Termikus, mechanikai és kémiai egyensúlyt feltételezve vizsgáljuk erősen kölcsönható anyag fajlagos ritkaságát. A ritkaságőrző és ritkaságváltoztató gyenge folyamatok egyensúlyán keresztül az egyes ízek relatív súlya erősen függ a leptontúlsúlytól, melyet ma még nagyon kevéssé számszerűsíthetünk. Ha a Világegyetemben a leptonok (antileptonok) vannak túlsúlyban, mégpedig nagyságrendileg annyira, mint a barionok, akkor a ritka kvarkok el vannak nyomva (felszaporodnak). A hadronfázisban a ritka barionok száma felmehet a barionszám feléig.

1. Introduction

Witten's (1984) suggestion of a baryon concentration process in the confinement transition of the early universe has stimulated many activities. Alcock et al. (1987) and Fuller et al. (1988) consider the baryon concentration in the last quark islands at the end of the confinement transition as result of an extremely low transmission rate of baryonic charge from the quark matter regions into the hadron matter regions. In fact, in their model some small islands with very low specific entropy ($s \approx 1$) immersed in normal matter with $s \approx 10^{10}$ survive up to nucleosynthesis. Needless to say that the outcome of nucleosynthesis in $s \approx 1$ regions drastically differs from the standard scenario. In addition, as Applegate, Hogan and Scherr (1987) point out, due to the different diffusion lengths of protons and neutrons, the baryon concentration is accompanied by strong inhomogeneities of protons and neutrons. Otherwise, as Witten (1984) states, "There is a rich element of wishful thinking here,...". Indeed, Applegate and Hogan (1985) and Kurki-Suonio (1988) remark that there are processes which tend to prevent the baryon concentration.

In his paper, Witten (1984) speculated that the highly baryon-enriched quark islands might transform into strange matter and might survive as dark matter. Farhi and Jaffe (1984) and Alcock and Farhi (1985), however, show that this turns out to be unlikely. Interesting in itself, even when speculating on the possible matter-stabilising role of strangeness, these authors do not include strangeness into the cosmic evolution of quark matter during the confinement transition, but only when the quark islands have been cooled down to nearly zero temperature.

It is the goal of the present paper to consider the evolution of the flavor composition of strongly interacting matter during the cosmic confinement transition. To have a generic framework for the above mentioned concentration process we restrict ourselves here onto full equilibrium, i.e. the evolution is considered as sequence of equilibrium situations. The confinement transition time scale of the early universe, $t \approx 10\ \mu sec$, supports this as first approximation.

Due to equilibrium with respect to strangeness-nonchanging ($\Delta S = 0$) and strangeness-changing ($\Delta S = 1$) weak interaction processes one finds for the chemical potentials of the down and strange quarks $\mu_d = \mu_s$. From this one expects a large fraction of strange matter in the early universe. This is the first difference of big bang and little bang (cf. Barz et al. 1988, Matsui et al. 1986, Kapusta and Mekjan 1986, Koch et al. 1986). The same situation one also meets in neutron stars with quark cores (Freedman and McLerran 1978).The second difference of big bang and little bang, and also to neutron stars, is that the lepton net numbers are conserved quantities. From this one expects that the actual flavor composition will depend mainly on the lepton excess.

Our paper is organised as follows. In section 2 we consider a simple toy model to clarify the order of magnitude of the chemical potentials. In section 3 we consider the phase diagramm and the flavor dynamics for up, down and strange quarks. Numerical results are discussed in section 4. The summary can be found in section 5. Some formulae used for the ideal gas approximation are listed in the appendix.

2. Prelude: Small chemical potentials

Here we present a simple toy model to clarify some basic features of the confinement transition. We discard isospin and strangeness for the moment being. As independent thermodynamical variables we take the nucleon chemical potential μ_N and the temperature T. Quark matter is modelled by an ideal gas of gluons and quarks (up and down), and the difference of the exact QCD and the perturbative vacua is parametrised by the bag constant B. In this approximation the pressure of the quark matter reads (cf. appendix)

$$p^Q = -B + {}^{37}/_{90}\, \pi^2 T^4 + \mu_q^{\,2} T^2. \qquad (2.1)$$

Hadron matter is mimicked by (cf. appendix)

$$p^H = {}^{9}/_{90}\, \pi^2 T^4 + b^N T \, (1 + {}^{1}/_{2}\, \mu_q^{\,2} T^{-2}) \qquad (2.2)$$

when regarding only for pions (taken for simplicity in the temperature range of interest in the ultrarelativistic approximation) and nucleons. The phase border line is constructed by exploiting the conditions of phase equilibrium, i.e. temperature, pressure and chemical equilibrium. The latter one reads $\mu_N = 3\,\mu_q$. With this we get in the regime of small chemical potentials

$$T_c(\mu_N) = T_o\ (1 - \mu_N^{\ 2}\ T_o^{\ -2}\ C), \qquad (2.3)$$

$$C = ({}^1/_9 - {}^1/_2\ \bar{b}_o^N)\ /\ (\ 4\ {}^{34}/_{90}\ \pi^2 - (2.5\ \ +m_N/T)\ \bar{b}_o^N),$$

where the critical temperature $T_o = T_c(\mu_N = 0)$ is defined by

$$-B + {}^{34}/_{90}\ \pi^2\ T_o^{\ 4} - b_o^{\ N}\ T_o = 0, \qquad (2.4)$$

$b_o = b(T_o)$. The phase borderline in the T-μ plane is a parabola (see fig.1a, lhs); inside is the confined region (world of hadrons) and outside the deconfined region (world of quarks and gluons). The corresponding T-n diagramm (n is the net baryon density) is displayed in fig.1a, rhs. The hatched area indicates the coexistence region. A state within the coexistence region is represented by an appropriately weighted superposion of states on the phase boundaries at given temperature. The phase border lines in the T-n diagramm are again parabolae (see fig.1a, rhs),

$$T_c(n^Q) = T_o\ (1 - (n^Q)^2\ \ T_o^{-3/2}\ \ \ {}^{81}/_4\ C), \qquad (2.5)$$

$$T_c(n^H) = T_o\ (1 - (n^H)^2\ \ b_o^{\ -2}\ C).$$

We note that, according to the present belief, the critical temperature is a fixed (but up to now uncertainly known) number in the order 100 - 300 MeV, while the bag constant is regarded as suitable parametrisation of the critical temperature. Therefore, B is a model dependent parameter. We will here discuss our results as function of T_o.

Witten (1984) observed that the ratio of the baryon excess densities in the quark and hadron phases,

$$n_Q/n_H = {}^2/_9\ \ T_o^{\ 3}\ b_o^{\ N}, \qquad (2.6)$$

is large for not too high critical temperatures (180, 14, 5, 3 for T_o = 100, 150, 200, 250 MeV, respectively). Bonometto et al. (1986) and Bonometto and Masiero (1986) considered this ratio in a model of many hadrons which interact via a strongly repulsive hard core; they found even more enlarged values, also at high temperatures. However, the predictive power of such models is rather uncertain as long as lattice QCD results cannot support them by direct calculations with finite chemical potentials.

As long as dissipative effects in the expanding universe can be neglected, the cosmic expansion proceeds isentropic, i.e the specific entropy $s = \sigma/n$ remains constant. (In case of vanishing baryon density the comoving entropy density $\sigma * R^3$ is conserved.) Present observations point to $s \approx 10^{10}$. This fixes the path of the universe in the thermodynamical state space (see fig. 1a). The entropy density σ is the superposition of background and strongly interacting components. The background entropy density is dominated by ultrarelativistic particles ($\gamma, e, \mu, \nu_e, \nu_\mu, \nu_\tau$, cf. Kämpfer and Schulz 1984),

$$\sigma_{bg} = 4\ a_{bg}\ T^4\ , \qquad a_{bg} = 14.25\ \pi^2/90. \tag{2.7}$$

In the quark phase we have therefore (cf. appendix)

$$T = s\ \mu_N\ {}^2/_9\ \ a_Q^{-1}, \tag{2.8}$$

$$a_Q = 4\ a_{bg} + {}^{75}/_{45}\ \pi^2,$$

while in the hadron phase

$$T = s\ \mu_H\ /\ a_H, \tag{2.9}$$

$$a_H = (4\ a_{bg} + {}^4/_{30}\ \pi^2)\ \tilde{b}^N + 2.5 + m_N\ /\ T.$$

Therefore, in the quark phase μ/T = const, while in the hadron phase μ/T is strongly increasing with decreasing temperature (see fig.2). It is evident that in a $s \approx 10^{10}$ universe, $\mu \ll T$ is an excellent approximation for $T > 50$ MeV. (Strickly speaking, below 150 MeV the ultrarelativistic approximation for pions and muons does not longer hold).

The isentropic line (2.7) intersects the phase border (2.3) at

$$T_Q^{\ *} = T_o\ (1 - s^{-2}\ a_Q^{\ 2}\ \ {}^{81}/_4), \tag{2.10}$$

$$\mu_{NQ}^{\ *} = T_o\ s^{-1}\ \ {}^9/_2\ \ a_Q,$$

and the hadronic isentropic line (2.8) intersects the border at

$$T_H^{\ *} = T_o\ (1 - s^{-2}\ C\ a_H^{\ 2}), \tag{2.10'}$$

$$\mu_{NH}^{\ *} = T_o\ s^{-1}\ a_H.$$

From these equations one derives that the temperature change during the adiabatic phase transition is negligible, $T_H^{\ *}/\ T_Q^{\ *} = O(s^{-2})$, while the change in the chemical potential can be considerable (see fig.2)

$$\mu_Q^{\ *}/\mu_H^{\ *} = {}^9/_2\ \ a_Q\ /\ a_H. \tag{2.11}$$

Interesting in itself, for $T < 250$ MeV this ratio is smaller than unity, thus indicating that the chemical potential increases during the adiabatic phase transition. Accordingly the temperature decreases. This is contrary to the general expectation (cf. Barz et al. 1984 for the little bang at large baryon density) that the released latent heat causes a reheating. Only for $T_o > 250$ MeV we find in the present model a reheating accompanied by a decrease of the chemical potential. Since the critical temperature depends on the vacuum energy density, it is clear that in case of smaller values of T_o or B the vacuum energy density is too small, for overcompensating the cooling, by released latent heat. Contrary to the reheating scenario, in the cooling scenario the densities of both the quarks and the hadrons grow during the hadronisation (see fig. 1b). This density increase might be considerable (e.g. a factor 100 for T_o = 100 MeV). But in equilibrium the density ratio of both phases remains constant according to eq. (2.6); it does not mean that one of the phases accumulates the baryons of the other one. Since the weight of the denser phase continuously decreases, the net effect is clearly the dilution of the total baryon density (see fig.1b).

3. Flavor dynamics

Now we extend the previous toy model and include explicitly up, down and strange quarks. In the hadron world we include neutrons (N = udd), protons (P = uud), kaons ($K^o = d\bar{s}$, $K^+ = u\bar{d}$), lambdas (Λ = uds) and sigmas (Σ^- = dds, Σ^+ = uus). The pions play in the quark picture of hadrons a particular role: sometimes they are regarded as elementary Goldstone modes and have not so a simple quark composition (Kocic 1988); otherwise in the naive quark model (Close 1979) they are regarded as simple quark composites ($\pi^o = u\bar{u}$ + some admixture of $d\bar{d}$, $\pi^+ = u\bar{d}$). The corresponding antiparticles are obtained by conjugating the quark content (but for π^o). The abundancies of heavier particles such as c, t and b quarks, ρ, ω, η mesons and Σ^o, Δ, Ω, Ξ baryons in the temperature range of interest are suppressed by the Boltzmann factors exp(-m/T). So for our goal, and also to maintain the transparency of the model, we include only the lightest representatives of the particle multiplets. We also do not include the double and triple strangeness charged baryons. Their inclusion is straightforward, but of minor effect.

Chemical equilibrium implies the relations of the chemical potentials

$$\mu_N = \mu_u + 2\mu_d,\ \mu_P = 2\mu_u + \mu_d,\ \mu_{Ko} = \mu_d - \mu_s,\ \mu_{K+} = \mu_u - \mu_s,$$
$$\mu_\Lambda = \mu_u + \mu_d + \mu_s,\ \mu_{\Sigma-} = 2\mu_d + \mu_s,\ \mu_{\Sigma+} = 2\mu_u + \mu_s,$$
$$\mu_{\pi o} = 0,\ \mu_{\pi+} = \delta_\pi\,(\mu_u - \mu_d), \tag{3.1}$$

where $\delta_\pi = 1$ means that we take the pions with naive quark contents, while $\delta_\pi = 0$ means the pions are regarded as Goldstone particles. Thus we end up with the independent chemical potentials μ_u, μ_d and μ_s and the temperature T as state variables.

To construct the phase boundary in the $T-\mu_u-\mu_d-\mu_s$ state space we have to specify the quark masses. Up and down quarks are taken as usual as ultrarelativistic particles. Koch, Müller and Rafelski (1986) and also Gasser and Leutwyler (1981) advocate a strange quark mass of $m_s \approx 150$ MeV. Otherwise the naive quark model (Close 1979) points to larger values, $m_s \approx 300 - 500$ MeV. In the former case the ultrarelativistic approximation for $T > 150$ MeV is appropriate, while in the latter case the Boltzmann approximation should be taken. We include both possibilities.

Thermal, mechanical (i.e. pressure) and chemical equilibrium gives for the phase boundary, when using the expressions for the pressure as stated in the appendix,

$$T_c = T_o(1 - a[\mu_u^2 - \mu_d^2] - c\,\mu_s^2 - d\,\mu_u\,\mu_d - e[\mu_u\,\mu_s - \mu_d\,\mu_s]), \qquad (3.2)$$

$$a = (-{}^1/_2 + {}^1/_6\,\delta_\pi + {}^5/_4\,\bar{b}_o^N + {}^1/_8\,\bar{b}_o^K + {}^1/_4\,\bar{b}_o^\Lambda + \bar{b}_o^\Sigma)/N,$$

$$c = (-{}^1/_2\,\bar{\beta}_o^s + {}^1/_4\,\bar{b}_o^K + {}^1/_4\,\bar{b}_o^\Lambda + {}^1/_2\,\bar{b}_o^\Sigma)/N,$$

$$d = (-{}^1/_3\,\delta_\pi + 2\,\bar{b}_o^N + {}^1/_2\,\bar{b}_o^\Lambda)/N,$$

$$e = (-{}^1/_4\,\bar{b}_o^K + \bar{b}_o^\Sigma + {}^1/_2\,\bar{b}_o^\Lambda)/N,$$

$$N = -2\,{}^{68}/_{90}\,\pi^2 + \beta_o^s\,(\delta^s - 4) + ({}^5/_2 + m_N/T_o)\,\bar{b}_o^N + {}^1/_2\,({}^5/_2 + m_K/T_o)\,\bar{b}_o^K + {}^1/_2\,({}^5/_2 + m_\Lambda/T_o)\,\bar{b}_o^\Lambda + ({}^5/_2 + m_\Sigma/T_o)\,\bar{b}_o^\Sigma,$$

where we used the abbreviations

for light strange quarks: $\beta^s = {}^1/_2\,{}^{21}/_{90}\,\pi^2$, $\bar{\beta}^s = 1$, $\delta^s = 0$,

for heavy strange quarks: $\bar{\beta}^s = \beta^s = {}^3/_2\,\bar{b}^s$, $\delta^s = {}^3/_2 - m_s/T$,

and the bag constant is determined by the critical temperature via

$$B = T_o^4\,({}^1/_2\,{}^{68}/_{90}\,\pi^2 + \beta_o^s) - b_o^N\,T_o - b_o^K\,T_o/2 - b_o^\Lambda\,T_o/2 + b_o^\Sigma\,T_o. \qquad (3.3)$$

In constrast to the little bang we have to impose in the big bang the equilibrium with respect to the weak interactions

$$d \leftrightarrow u + e^- + \bar{\nu}_e \quad \text{and} \quad s \leftrightarrow u + e^- + \bar{\nu}_e \tag{3.4}$$

which imply

$$\mu_d - \mu_u - \mu_e + \mu_{\nu e} = 0,$$
$$\mu_s - \mu_u - \mu_e + \mu_{\nu e} = 0. \tag{3.4'}$$

A consequence of the equilibrium with respect to these weak interactions is the equality of the down and strange chemical potentials, $\mu_s = \mu_d$ (cf. Freedman and McLerran 1978). The corresponding muonic reactions are

$$d \leftrightarrow u + \mu^- + \bar{\nu}_\mu \quad \text{and} \quad s \leftrightarrow u + \mu^- + \bar{\nu}_\mu \tag{3.5}$$

which in turn imply

$$\mu_d - \mu_u - \mu_\mu + \mu_{\nu\mu} = 0, \tag{3.5'}$$
$$\mu_s - \mu_u - \mu_\mu + \mu_{\nu\mu} = 0.$$

The reaction equations (3.5) and (3.6) combine, e.g. to the pure leptonic weak reaction $e^- + \bar{\nu}_e \leftrightarrow \mu^- + \bar{\nu}_\mu$ (cf. Weinberg 1972). The possible hadronic weak reactions follow from eqs.(3.4), (3.5) and the above-given quark contents of the hadrons.

The particular equilibrium condition $\mu_s = \mu_d$ implies that the ratio of net strange quark density to net down quark density in the quark phase is

$$n_s/n_d = \bar{\beta}^s. \tag{3.6}$$

Thus the early universe, in the quark phase can be heavily charged up with strange quarks.

The intersection of the phase boundary (3.2) with the hypersurface $\mu_d = \mu_s$ describes the phase boundary in the $T-\mu_u-\mu_d$ state space

$$T_c = T_o\,(1 - \bar{a}\,\mu_u^2 - \bar{c}\,\mu_d^2 - \bar{d}\,\mu_u\,\mu_d)/N, \tag{3.7}$$
$$\bar{a} = -{}^1/_2 + {}^1/_6\,\delta_\pi + {}^5/_4\,\bar{b}_o^N + {}^1/_8\,\bar{b}_o^K + {}^1/_4\,\bar{b}_o^\Lambda + \bar{b}_o^\Sigma,$$
$$\bar{c} = -{}^1/_2\,[1 + \bar{\beta}_o^s] + {}^1/_6\,\delta_\pi + {}^5/_4\,\bar{b}_o^N + {}^1/_8\,\bar{b}_o^K + \bar{b}_o^\Lambda + {}^5/_2\,\bar{b}_o^\Sigma,$$
$$\bar{d} = -{}^1/_3\,\delta_\pi + 2\,\bar{b}_o^N - {}^1/_4\,\bar{b}_o^K + \bar{b}_o^\Lambda + \bar{b}_o^\Sigma.$$

This surface is schematically displayed in fig.3. Similar to the simple toy model of the previous section, here the temperature and the chemical potentials change in general during the phase transition.

The actual path in the thermodynamical state space is restricted by the conservation laws:

(i) conservation of electric charge:

$$n_Q = 0, \qquad (3.8)$$
$$n_Q = -n_e - n_\mu + x\,[2\,n_u - n_d - n_s] + (1-x)\,[n_p + n_{K+} + n_{\Sigma+} - n_{\Sigma-} + n_{\pi+}]$$

(here we regarded the charges of u,d,s quarks with $^2/_3$, $-^1/_3$, $-^1/_3$ and, by definition, $3n_{u,d,s}$ are the particle net numbers, while $n_{u,d,s}$ are the baryon net numbers; the other n_i are the net numbers, i.e. the particle minus antiparticle numbers; in particular the K^- and π^- are regarded as antiparticles to K^+ and π^+);

(ii) conservation of baryon number:

$$R^3\,n_B = \text{const}, \qquad (3.9)$$
$$n_B = x\,[n_u + n_d + n_s] + (1-x)\,[n_p + n_N + n_\Lambda + n_{\Sigma-} + n_{\Sigma+}]$$

(R is the scale factor of the Robertson Walker metric, not needed explicitly in what follows);

(iii) conservation of electron lepton number:

$$R^3\,n_E = \text{const}, \qquad (3.10)$$
$$n_E = n_e + n_{\nu e};$$

(iv) conservation of muon lepton number:

$$R^3\,n_M = \text{const}, \qquad (3.11)$$
$$n_M = n_\mu + n_{\nu\mu};$$

(v) conservation of tau lepton number:

$$R^3\,n_T = \text{const}, \qquad (3.12)$$
$$n_T = n_{\nu\tau}$$

(in the temperature range of interest the taus have annihilated or decayed; so the tau lepton charge is carried by the tau neutrinos).

In eqs. (3.8) - (3.12) we include the possibility that the volume fraction x stays in the quark phase, while the fraction 1-x is in the hadron phase. x = 1 and 0 correspond to pure quark and hadron phases, respectively. Since the conservation laws for the baryon and lepton charges have the same form we can write

$$n_E = \alpha_E\,n_B, \quad n_M = \alpha_M\,n_B, \quad n_T = \alpha_T\,n_B, \qquad (3.13)$$

where the $\alpha_{E,M,T}$ are constants. Unfortunately one knows nothing on the present densities of the various neutrinos. Therefore an often used *ad hoc* choice is $\alpha_{E,M,T} = 0$. This would mean that either the net number of every lepton sort, including the neutrinos, vanis ($\mu_{e,\mu,\tau,\nu e,\nu\mu,\nu\tau} = 0$) or the number of massive leptons is balan by the corresponding antineutrinos. Otherwise the grand unif theories (GUT) suggest that a baryon excess can be generated simultaneously with a lepton excess in $\Delta(B - L) = 0$ processes during the GUT symmmetry breaking phase transition at $T \approx 10^{15}$ GeV. This would favor $\alpha_{E+M+T} = 1$ if one assumes a total symmetric universe at Planck temperature. In addition, there are also baryosynthesis processes in the electroweak symmetry breaking at $T \approx 100$ GeV via some anomalies. All such processes are still under debate. Therefore, the relation of the baryon excess to the lepton excess is poorly known. We will take in what follows the excess constants α as free parameters and discuss our results as function of them.

The four conservation laws (3.8)-(3.13) and the four equilibrium conditions (3.4)-(3.5') give partially degenerated restrictions onto the nine variables T, $\mu_{u,d,s,e,\mu,\nu e,\nu\mu,\mu\tau}$ which determine the path of the cosmic matter in the thermodynamic state space, e.g. in the form $\mu_u = \mu_u(T,\mu_d)$. In absence of dissipative effects the trajectory is then uniquely determined by the isentropic condition s = const, which can be represented, e.g. in the form $T = T(\mu_d)$. The explicit form of $\mu_u(\mu_d,T)$ reads

$$\mu_u = \mu_d A, \qquad A = Z / N \tag{3.14}$$

$$Z = 1 + x\,[\alpha + {}^{3}/_{2}]\,(1 + \bar{\beta}^{e}) + (1-x)\,[\bar{b}^{N}\,(2\,\alpha - 1)\ {}^{g}/_{4}$$
$$+ \bar{b}^{K}\ {}^{g}/_{4} + \bar{b}^{\Lambda}\, 3\,\alpha + \bar{b}^{\Sigma}\,({}^{g}/_{2} + 6\,\alpha) + {}^{3}/_{2}\,\delta_{\pi}],$$

$$N = 1 + x\,[3 - \alpha] + (1-x)\,[\bar{b}^{N}\,(1-\alpha)\ {}^{g}/_{2} + \bar{b}^{K}\ {}^{g}/_{8} - \bar{b}^{\Lambda}\ {}^{3}/_{2}\,\alpha$$
$$+ \bar{b}^{\Sigma}\,({}^{g}/_{2} - 3\,\alpha) + {}^{3}/_{2}\,\delta_{\pi}],$$

$$\alpha = \alpha_E + \alpha_M.$$

Note that α_T does not enter into the equations since the taus do not participate in the electric charge balance. Also due to the large tau mass, they do not participate in the equilibrium reactions with other light leptons; therefore the tau neutrino chemical potential does not couple to other chemical potentials. We also emphasise the necessity to include properly the muonic reactions (otherwise one would get a term $^{1}/_{2}$ instead 1 in front of the nominator and demominator of A and would imply $\alpha_M = 0$; but then one would discard the reactions $e^- + \bar{\nu}_e$ <-> $\mu^- + \bar{\nu}_\mu$ and would arbitrarily set $\mu_\mu = 0$).

The isentropic lines are determined by

$$s\,\mu_d / T = Z_1 / N_1, \tag{3.15}$$

$$Z_1 = 4\,a_{bg} + x\,[{}^{74}/_{45}\,\pi^2 + \bar{\beta}^{s}] + (1-x)\,[{}^{4}/_{80}\,\pi^2 +$$
$$\bar{b}^N\,({}^{5}/_{2} + m_N/T) + \bar{b}^K\,({}^{5}/_{2} + m_K/T)/2 + \bar{b}^\Lambda\,({}^{5}/_{2} + m_\Lambda/T)/2$$
$$+\bar{b}^\Sigma\,({}^{5}/_{2} + m_\Sigma/T)/2],$$

$$N_1 = x\,[A + 1 + \bar{\beta}^{s}] + (1-x)\,[\bar{b}^N\,(1 + A)\,{}^{3}/_{2} + \bar{b}^\Lambda\,(A + 2)/2$$
$$+\,\bar{b}^\Sigma\,(A + 2)],$$

for $x = 1$ and 0. The ratio $s\mu_d/T$ is, of course, very similar to $s\mu_N/T$ (see fig.2). Eqs. (3.15) with (3.14) entirely fix the trajectory of the cosmic matter in the thermodynamic state space.

4. Discussion of the numerical results

Having the relations of the quark chemical potentials, eq.(3.14) and $\mu_d = \mu_s$, at disposal we can calculate the various ratios of particles net numbers during the confinement transition. Obviously these ratios depend on the phase weight x. In particular we are interested in the ratios at the beginning (x=1) and at the end (x=0) of the transition. It turns out that all ratios considered below are monotonous functions of x. Therefore, we present results only for x=1, 0. We consider two representative critical temperatures T_o = 150 and 250 MeV. Also we study the influence of the strange quark mass.

The ratio of up to down quarks is

$$n_u/n_d = A \tag{4.1}$$

which is displayed in fig.4. O[illegible]serves for values of α smaller than -1 that the up quarks are suppressed or up antiquarks are in excess of the up quarks. At $\alpha \approx 0$ up and down quark abundancies are fairly balanced, and at higher values of α the up quarks start to dominate. The influence of the strange quark mass is rather moderate. At lower critical temperature the up-down asymmetry diminishes during the phase transition. At high transition temperature the up-down asymmetry increases slightly, when pions are taken as Goldstone modes (i.e. without definite quark contents). For higher transition temprratures our model predicts stronger asymmetry effects which depend sensitively on the value of α.

The ratio

$$n^Q_s \,/\, n^Q_{no} = n_s \,/\, (n_u + n_d) = \bar{\beta}^s \,/\, (1 + A) \qquad (4.2)$$

describes the ratio of net strange quark number to the u + d net number (see fig.5). For heavy strange quarks the strange quark fraction is obviously suppressed (rhs of fig.5). In a universe with $\alpha \approx -2$ the strange quarks are as frequent as u and d quarks together. Thus strange quarks carry a half of the baryon charge. In a universe with $\alpha \approx 2$ the strange quarks are surprisingly suppressed.

On fig.6 the ratio

$$n^H_{ch}/n^H_{ne} = \frac{n_p + n_{\Sigma+} + n_{\Sigma-}}{n_N + n_\Lambda} = \frac{\bar{b}^N(2A + 1) + \bar{b}^\Sigma(2A + 4)}{\bar{b}^N(A + 2) + \bar{b}^\Lambda(A + 2)} \qquad (4.3)$$

of charged to neutral baryons is displayed. One observes a rather balanced distribution of charged and neutral baryons (this corresponds to the proton-neutron symmetry often assumed).

The ratio of net strangeness in the quark regions to the net strangeness in the hadron regions

$$\tilde{n}^Q_s/\tilde{n}^H_s = 3n_s/(n_\Lambda + n_{\Sigma+} + n_{\Sigma-} + n_{KO} + n_{K+}) \qquad (4.4)$$
$$= \bar{\beta}^s \,/\, [\bar{b}^\Lambda (A+1)/2 + \bar{b}^\Sigma (A + 2) - \bar{b}^K (A - 1)/4]$$

(see fig.7) clearly resembles somewhat to the baryon excess. In particular, for small transition temperature there is much more strangeness in the quark phase than in the hadron phase. This effect is due to the suppression of the abundancies of particles with mass $m > T$. At higher transition temperature this difference is strongly reduced. Note that in particular for $\alpha \approx 1$ there is an extreme surplus of strangeness in the quark phase, supposed the critical temperature is small enough. For higher transition temperature and $\alpha > 1$ this effect is reversed: the strangeness in the quark phase can be strongly suppressed.

On fig.8 the ratio of strangeness per baryon in the quark phase to the hadron phase,

$$n^Q_s/n^H_s = n_s/(n_\Lambda + n_{\Sigma+} + n_{\Sigma-})$$
$$= {}^2/_3\, \bar{\beta}^s/[\bar{b}^\Lambda(2+A)+\bar{b}^\Sigma(4+2A)], \qquad (4.5)$$

is displayed. This ratio resembles directly to the baryon excess. Accordingly, at small critical temperature the number of strangeness in the quark phase is much higher (a factor of ten) than in the hadron phase. At higher transition temperature and for heavier strange quarks, the strangeness is sheared between both phases.

Finally we mention that we have also studied the strange distribution in the hadron phase. The ratio of strange baryons to non-strange baryons,

$$n_s^H / n_{ns}^H = (n_\Lambda + n_{\Sigma+} + n_{\Sigma-})/(n_N + n_P) \qquad (4.6)$$
$$= [\bar{b}^\Lambda (A + 2) + \bar{b}^\Sigma (4 + 2A)] / [3 \bar{b}^N (A + 1)],$$

is displayed in fig.9. One observes abundancies of the strange baryons in the range 25 - 50%. Thus also in the hadron phase a considerable fraction of baryonic charge is carried by the strange particles.

5. Summary

In the present paper we study the flavor dynamics during the cosmic confinement transition exploiting equilibrium conditions. In particular we include equilibrium with respect to weak interactions. By this the hadronisation in the big bang differs mainly from the little bang hadronisation. The electric charge neutrality as well as baryon and lepton conservation restrict the flavor composition of strongly interacting matter. Unfortunately, the lepton excess numbers are presently poorly known, and some particle densities depend senitively on them. If the antilepton excess is in the same order of magnitude as the baryon excess, than a half of the baryon excess is carried by strange quarks. If the lepton excess is in the same order as the baryon excess (as predicted by GUT) then we find strongly suppressed strange quark abundancies. In some regions of the parameter space which we consider (such as critical temperature, lepton excess, strange quark mass) the quark phase can carry an extremely large net strangeness as compared with the hadron phase. But the ratio of strange baryons in the quark and hadron phase goes parallel to the general baryon excess, unless for heavy strange quarks and large transition temperature. In the hadron phase up to 50% of the baryon excess is carried by strange baryons.

Finally we comment on the applicability of the equilibrium scenario. Our previous study (Kämpfer 1988) on supercooling effects within the framework of the nucleation theory demonstrates that for surface tensions $\sigma_{surf} = T^3 \sigma_o$ with $\sigma_o < 3.5$ the universe reheats, after a supercooling period, rapidly to T_o. Therefore for not too large values of σ_o our present equilibrium scenario is applicable.

Acknowledgement

Enlightening discussions with Drs. B.Lukacs and M.I.Gorenstein are gratefully acknowledged.

Appendix: Approximations for ideal gases

Thermodynamical properties of an ideal gas can be calculated from the thermodynamical potential $p(T,\pm\mu)$

$$p(T,\pm\mu) = g/(6\pi^2) \int dp\, p^4/\varepsilon \quad [\exp\{(\varepsilon \pm \mu)/T\} + I]^{-1}, \qquad (A.1)$$

$$\varepsilon = (m^2 + p^2)^{1/2}$$

via

$$n(T,\pm\mu) = \partial p(T,\pm\mu)/\partial\mu, \qquad s(T,\pm\mu) = \partial p(T,\pm\mu)/\partial T, \qquad ((A.2)$$

$$\mu\, n = e + p - T\, s.$$

As usual, $I = \pm 1$ for Fermions and Bosons. We neglect Bose condensate states, what is justified in the thermodynamical limit as long as $|\mu| < m$ for the respective particles. Since we are dealing with small chemical potentials, $\mu/T \approx 10^{-10}$, this is correct for pions and kaons. In case of the gluons the same arguments as for photons apply. In the case that antiparticles to the respective particles exist (such as for all Fermions and π^+, K^o, K^+), their pressure is $p(T,\mu) = p(T,+\mu) + p(T,-\mu)$, and the net number density reads correspondingly $n(T,\mu) = n(T,+\mu) - n(T,-\mu)$ (Take care for quarks! We reserve for them n as baryon net number being a third of the total net number), while for particles which possess no antiparticles (such as gluons, photons, π^o), $p(T,\mu) = p(T,\mu{=}0)$.

The following Fermions are taken in ultrarelativistic approximation:

u,d,s quarks ($g = 6$), electrons ($g = 2$), electron, muon, tau neutrinos ($g = 1$):

$$p = g\, [\,^7/_8\, ^2/_{90}\, \pi^2 T^4 + \,^1/_{12}\, \mu^2 T^2 + \mu^4 / (24\pi^2)] \qquad (A.3)$$

while the Boltzmann approximation is used for s quarks ($g = 6$), nucleons (protons, neutrons; $g = 2$), lambdas ($g = 2$), sigmas (Σ^+,

Σ^-; $g = 2$):

$$p = {}^1/_4\ g\ b\ T\ (1 + {}^1/_2\ (\mu/T)^2 + \ldots), \qquad (A.4)$$

$$b = (2\ m\ T\ /\ \pi)^{3/2}\ \exp(-m\ /\ T), \qquad \bar{b} = b\ T^{-3},$$

where the pressure already includes the antiparticle contributions and g means the degeneracy without antiparticles.

Ultrarelativistic Bosons without antiparticles (i.e. $\mu = 0$) are photons ($g = 2$), gluons ($g = 8$) and π^0 ($g = 1$):

$$p = {}^1/_{30}\ g\ \pi^2\ T^4 \qquad (A.5)$$

and also the charged pions which are regarded either as mutually conjugate particles ($\mu_{\pi+} = -\mu_{\pi-}$, $g = 1$)

$$p = g\ [{}^2/_{90}\ \pi^2\ T^4 + {}^1/_6\ \mu^2\ T^2 + \ldots] \qquad (A.6)$$

or as Goldstone modes ($\mu_{\pi\pm} = 0$)

$$p = {}^2/_{90}\ \pi^2\ T^4. \qquad (A.7)$$

The kaons also being Bosons, are taken in the Boltzmann approximation (A.4) with $g = 1$ for K^+, K^0 and $\mu_{K+} = -\mu_{K-}, \mu_{K0} = -\mu_{\bar{K}0}$.

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Figure captions

Fig.1a. Schematic plot of the phase diagramm. T_0 is the critical temperature. In the $T-\mu$ state space (lhs) the heavy line is the phase boundary, while in the T-n space (rhs) the hatched area is the coexistence region; left to the phase boundaries is the hadron world with confinement, and right to the boundaries is the world of deconfined quarks and gluons. Depicted are also isentropic lines: 1 - the latent heat is sufficiently large to ensure an reheating, 2 - the transformation is accompanied by cooling.

Fig.1b. Schematic spatial plot of the hadron bubbles (H) in the quark phase (Q) at the beginning of the transformations (lhs) and of quark bubbles in hadron matter at the end of the transformation (rhs). During the transformation the density jump $\Delta = n^Q/n^H$ remains constant according to eq. (2.6) since the temperature change is in the order s^{-2}. In the reheating scenario 1 all densities decrease, and the net density change δ is larger than Δ. In the cooling scenario 2 all densities increase and one has $\delta < \Delta$.

Fig.2. The quantity $s\mu/T$ as function of T (heavy line: hadron matter + background, dotted line: quark matter + background). The dashed line (rhs scale) shows the ratio of chemical potentials before and after an isentropic transformation, μ^*_Q/μ^*_H, as function of the critical temperature T_0.

Fig.3. Schematic plot of the phase diagram in the $T-\mu_u-\mu_d$ state space. Encovered by the phase boundary (with some isotherms on it) is the hadron world, and outside is the world of quarks and

gluons. The heavy line depicts a possible path of the universe before, during (here with small reheating) and after (dotted section) the transformation.

Fig.4. Ratio of net up quark density to net down quark density as function of $\alpha = \alpha_E + \alpha_M$ for two different critical temperatures (upper part: T_o = 150 MeV, lower part: T_o = 250 MeV) and for light (lhs) and heavy (rhs) strange quarks.

Heavy lines: at the beginning of the transformation (x=1); dashed lines (pions are handled according to their naive quark contents) and dotted lines (pions are taken as Goldstone modes with vanishing chemical potential): at the end of the transformation (x=0).

Fig.5. The same as in fig.4, but for the ratio of strange quarks to the non-strange quarks.

Fig.6. The same as in fig.4, but for the ratio of charged to neutral baryons in the hadron phase. The strange quark mass effect turns out to be neligible.

Fig.7. The same as in fig.4, but for the ratio of net strangeness in the quark to the hadron phase.

Fig.8. The same as in fig.4, but for the ratio of strangeness per baryon in the quark to the hadron phase.

Fig.9. The same as fig.4, but for the ratio of the strange baryons to non-strange baryons in the hadron phase. The strange quark mass effect is negligible.

Figure 1a

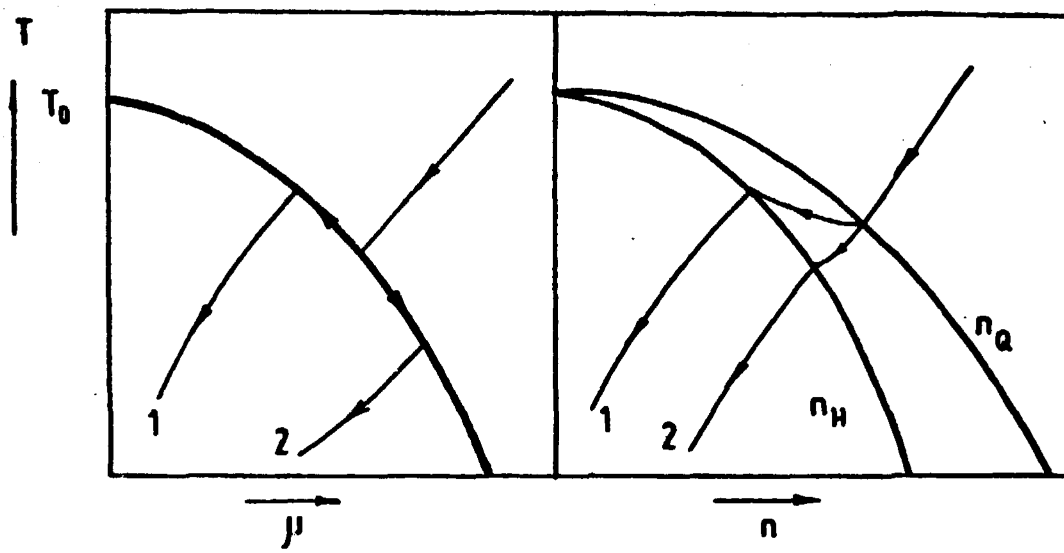
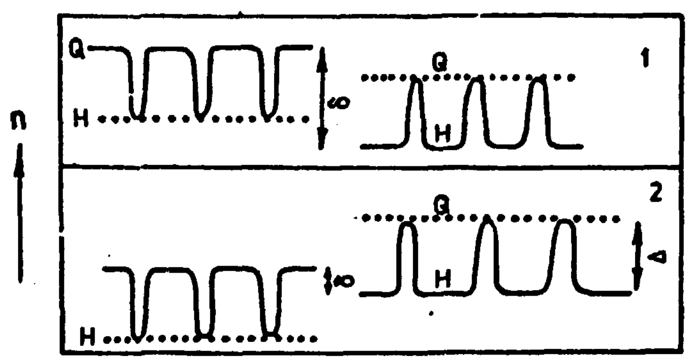


Figure 1b



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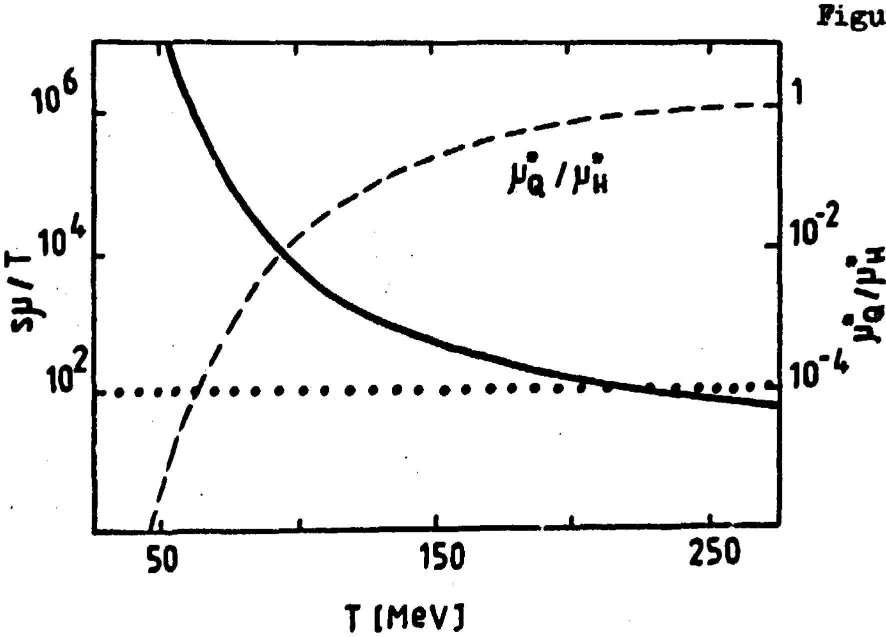

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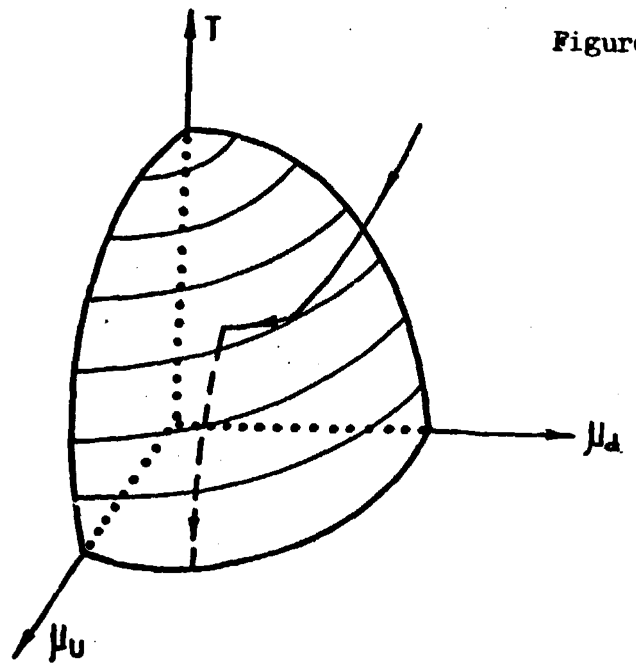

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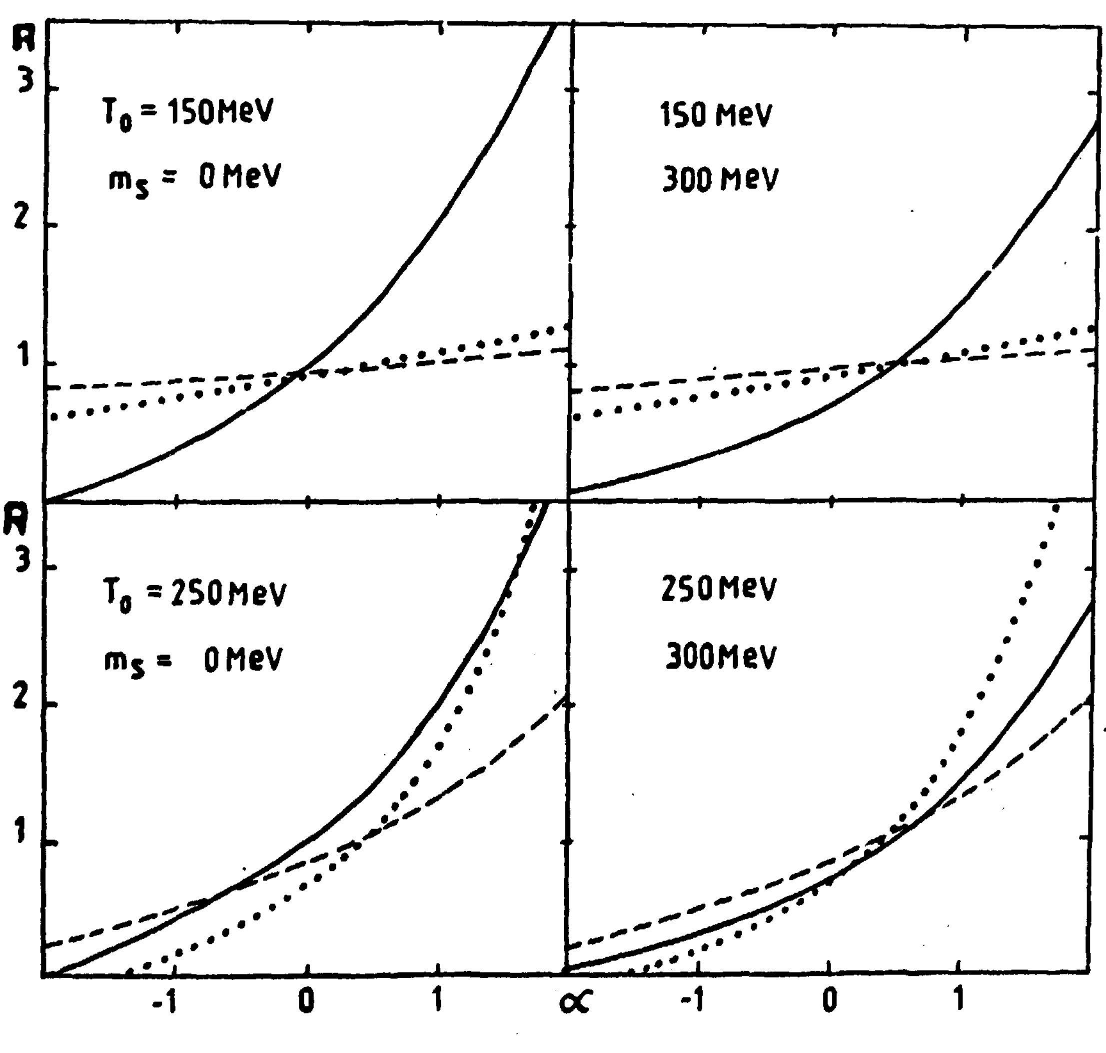
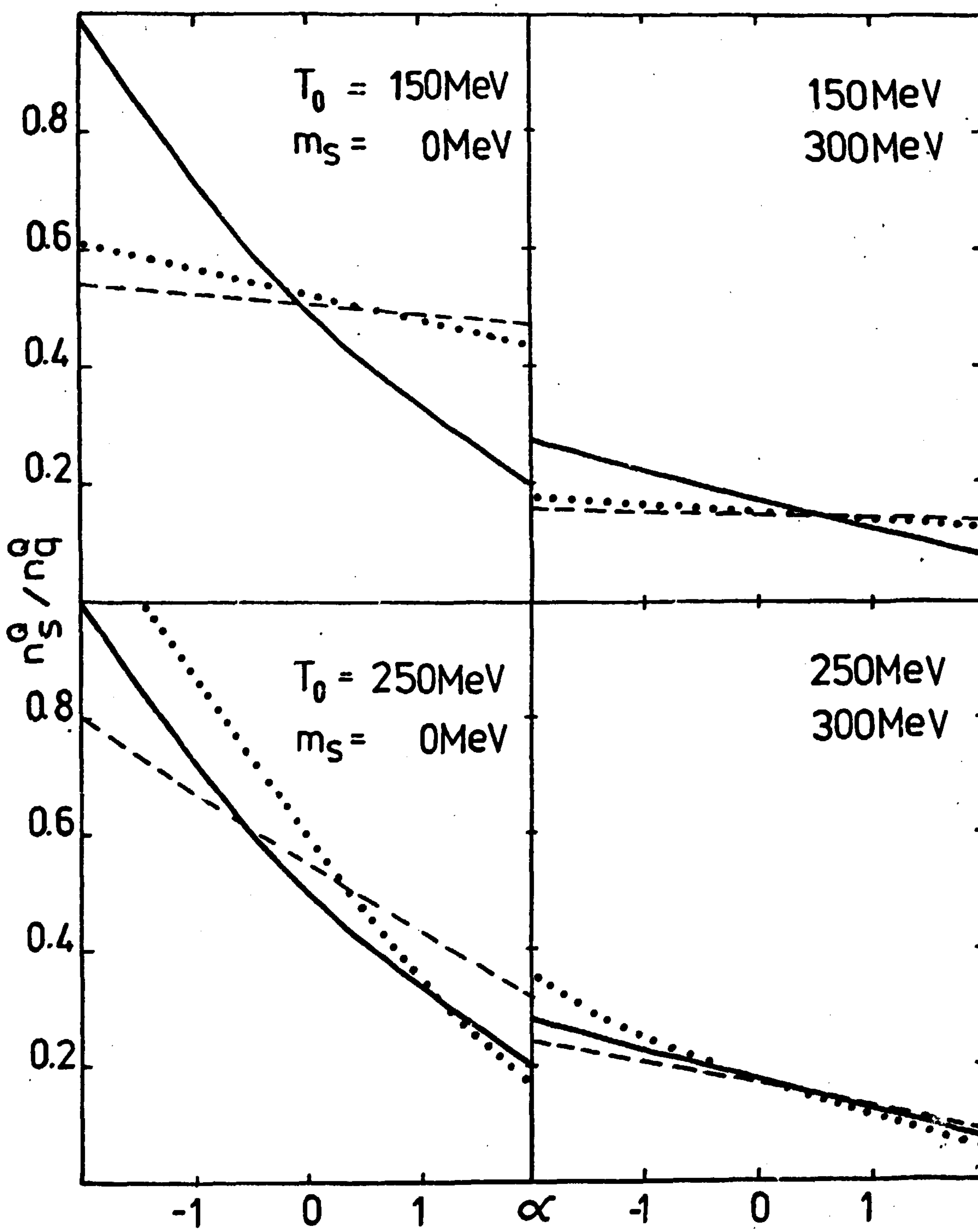


Figure 5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150MeV
300MeV
T0 = 250MeV
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250MeV
300MeV
nsa/nqa
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0.6
0.4
0.2
-1
0
1
α

Figure 6

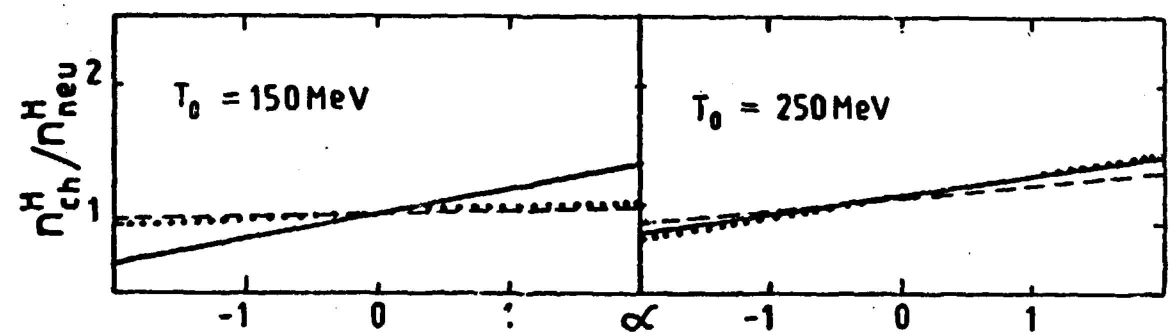


Figure 7

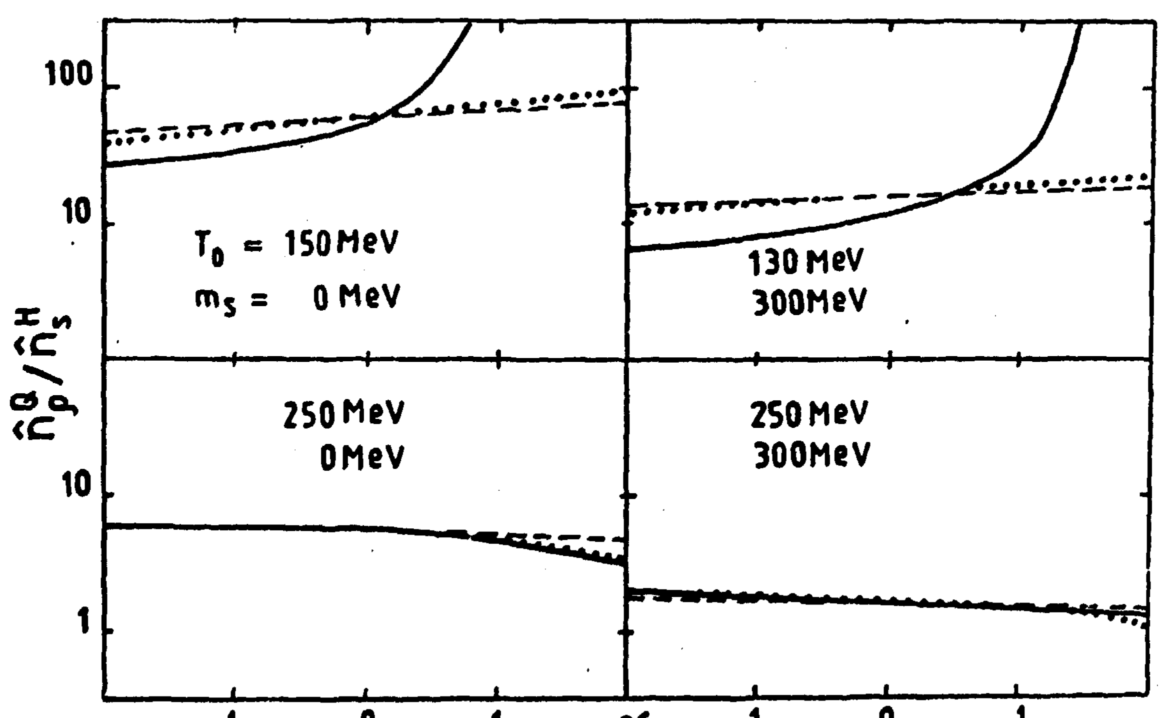


Figure 8

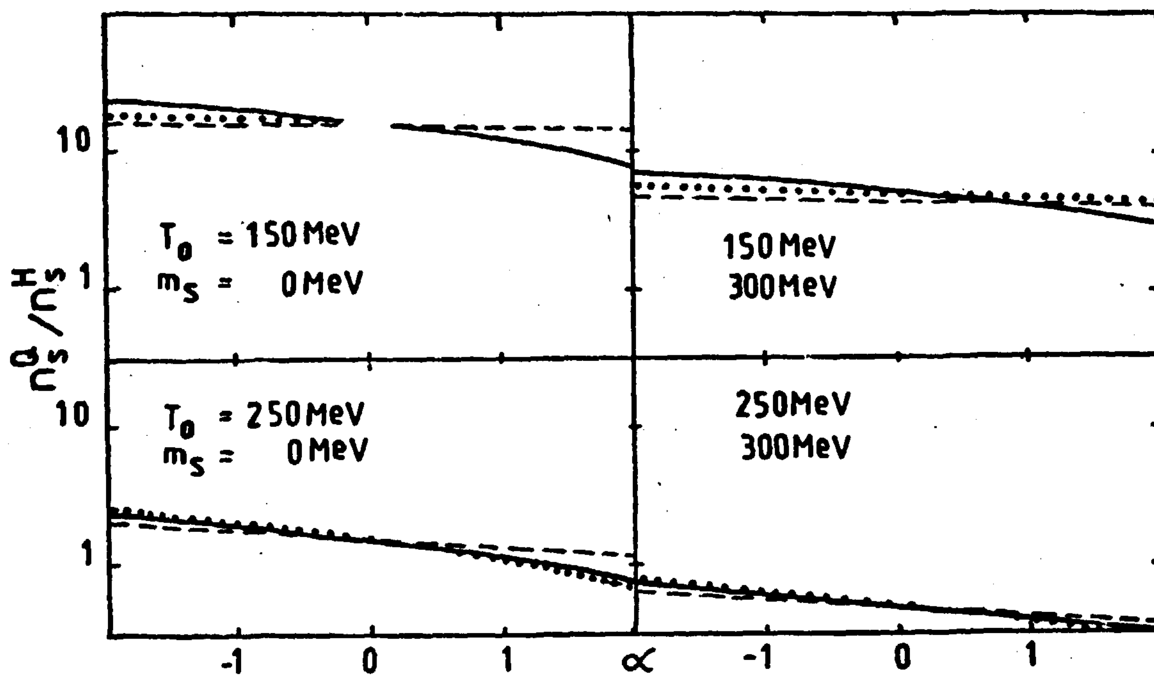
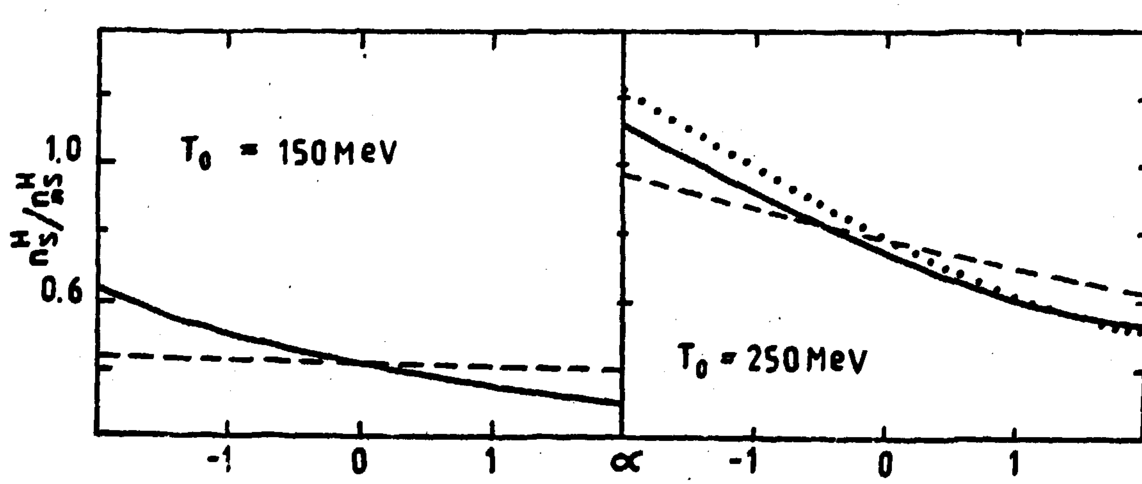


Figure 9



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