Coagulation of the quark-gluon plasma in $n\bar{n}$ annihilation and heavy ion collisions

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The consequences of the formation of a quark-gluon plasma and its later hadronization are investigated. It is assumed that the plasma is formed beyond a specified energy density in the two reactions considered, the total annihilation of antinucleons on nucleons and the central collision of heavy ions. A coalescence model for the formation of quark clusters or hadrons is formulated. The multiplicity and momentum distribution of pions and kaons from $n\bar{n}$ annihilation as also the ratio of the production of pions, kaons, and lambdas to protons are estimated for various collision energies of the heavy ions. The K^+ spectra at three angles are computed and compared with available experimental data.

I. INTRODUCTION

A question of great interest is whether a quark-gluon plasma could be created and recognized as such in laboratory experiments.[1] The nature of the transition from hadronic matter to the quark-gluon plasma is uncertain, but of more practical relevance is the difference between the two phases and the way the quark-gluon plasma hadronizes as it cools on expansion.

The most elementary argument for a quark-gluon phase derives from the phenomenological bag model[2] where the quarks are confined to within a nucleon radius, R. When nuclear matter is compressed so that the internucleon spacing becomes less than $2R$, the bags interpenetrate and their constituents form macroscopic objects and a gradual deconfinement to the quark-gluon phase[3] may occur. An internucleon spacing of less than $2R$ implies a nucleon density larger than $(3/4\pi R^3)$, and for a nucleon radius of 0.85 fm, a transition is expected to take place at about three times the normal nuclear density. Lattice quantum chromodynamics (QCD) calculations[4] support these naive arguments to some extent and indicate that a sharp deconfinement transition occurs[5] at 4—8 times normal nuclear density at low temperatures and at $T_c \sim 150$—250 MeV at low baryon densities.

The bootstrap model[6] for the hadronic phase predicts a transition to the quark-gluon phase when the energy density exceeds the bag pressure, $\epsilon \geq 4B$. With the value of $B^{1/4}=145$ MeV, as given by the MIT bag model,[7] the transition energy density is 0.25 GeV/fm^3, but a choice of $B^{1/4}=190$ MeV, consistent with the lattice QCD results, lifts this transition to 0.67 GeV/fm^3, close to the energy density within a hadron which is around 0.5 GeV/fm^3. On the other hand, evidence[8] exists that quarks in cold nuclei are probably at least partially deconfined and that the hopping of quarks from one nucleon to another could cause correlations and clustering that are observed in nuclei.

At high energy densities in the quark gas, the average momentum transfer in collisions is expected to be large and the asymptotic freedom of quark and gluon interactions may cause them to traverse the system freely, interacting with themselves. Gluons and quark-antiquark pairs are created in enough numbers to permit the system to be described by the statistical thermodynamics of a relativistic gas[9] of bosons and fermions with reasonable accuracy. At lower energy densities, such as near the transition point, perturbative QCD calculations would not be valid but the pressure and energy density can be fitted by lowest order perturbative expressions provided that one uses an *effective* QCD coupling constant, α_s, of 0.5—0.6 and a bag pressure of $B^{1/4}$ of between 150 and 190 MeV, treated as parameters. Nonperturbative interactions have been simulated by Monte Carlo techniques[10] and these suggest that the deconfinement is abrupt.

In the present work we consider two reactions where the energy density is fairly high, the annihilation of nucleons on antinucleons and the central collision of heavy ions. We *assume* that in each case a quark-gluon plasma is formed, an assumption which is reasonable in view of the rapid thermalization arising from the very strong quark-quark and quark-gluon interactions. The possible characteristic signals consequent to plasma formation, such as the production of strange particles particularly near the threshold energy are then investigated. We first discuss the model we employ in our calculations and later we present the results of this model for the two reactions considered.

II. THE QUARK-GLUON PLASMA

A. The formation of the plasma

The energy density ϵ in a plasma has been given[11] as a function of the bag constant B, the plasma temperature T ($\beta=1/T$), the chemical potential for the nonstrange

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quarks μ, and the QCD color coupling constant α_s:

$$\epsilon=B+\frac{8\pi^2}{15\beta^4}\left[1-\frac{15\alpha_s}{4\pi}\right]+\frac{7\pi^2}{10\beta^4}\left[1-\frac{50\alpha_s}{21\pi}\right]+\frac{1}{3\pi^2}\left[1-\frac{2\alpha_s}{\pi}\right]\left[\frac{\pi^2\mu^2}{\beta^2}+\frac{\mu^4}{18}\right]+\epsilon_s\ . \quad (1)$$

The contribution ϵ_s to the energy density from the strange quarks may be written as

$$\epsilon_s=\frac{6F}{\pi^2}\int q^2dq(q^2+m_s^2)^{1/2}\{\exp[\beta(q^2+m_s^2)^{1/2}]+1\}^{-1}\ , \quad (2)$$

where the mass of the strange quark is m_s and its momentum is q.

The form of Eq. (1) is correct even though a perturbative expansion may not be valid near the deconfinement transition, provided that B and α_s are treated as parameters. It is supposed that in the plasma the strange quark pairs are created through the processes[12] $gg\rightarrow s\bar{s}$ and $q\bar{q}\rightarrow s\bar{s}$. Within the small lifetime ($\sim 10^{-23}$ s) of the plasma, the strange quark density does not saturate except at extremely high temperatures. The approach to saturation in time has been computed,[13] but in the absence of a reliable estimate of the plasma lifetime the lack of saturation is defined by a fractional parameter, F, in Eq. (2). It varies from 0.1 to 0.2 for the reactions considered. The interactions of the strange quarks with themselves and with the other quarks and gluons have been neglected since this would cause only a small difference in an already small term.

In addition to the energy density expression of Eq. (1) the baryon number density in the plasma is

$$\rho_B=\frac{2}{3\pi^2\beta^3}\left[1-\frac{2\alpha_s}{\pi}\right]\left[\pi^2\mu+\mu^3\right]\ . \quad (3)$$

Finally the plasma will hadronize at an energy density where the bag pressure balances the pressure of the quarks and gluons as in the MIT bag model,

$$\epsilon=4B\ . \quad (4)$$

These four conditions given by Eqs. (1)—(4) with the input parameters B, α_s, and F determine the chemical potential for the nonstrange quarks as also the plasma temperature at the transition.

B. The hadronization process

The quark-gluon plasma must expand and cool back to the transition temperature before it nucleates into hadrons. At extremely high energies, perhaps above 60 GeV/fm^3 the quarks escape from the plasma and in the process lose their energy by the emission[14] of $q\bar{q}$ pairs (mesons) in a jet. The quark momenta in the plasma condition we consider here are, however, low, of the order of a few hundred MeV/c, and jet production is not possible. Instead we use a coalescence model to describe the hadronization near the transition point. Nucleon cluster formation[15,16] in heavy ion collisions has been explained in a similar picture, but with the criterion that nucleons with less than a specified relative momentum coalesce into a cluster. In our model we postulate that the probability of forming a hadron is proportional to the probability of finding three quarks (or a quark-antiquark pair) of appropriate color, spin, and flavor within a hadronic volume, typically a sphere of radius 0.8 fm. The great strength of the QCD color interaction motivates this assumption of an interaction volume and indeed a suggestion has been made[16] that this is a more appropriate picture even for nucleons. The appropriate combinations of quantum numbers in a three quark cluster form a color singlet baryon which can then escape from the plasma.

The probability, $W_B(p)$ of forming a baryon of four-momentum, p, is

$$W_B(p)=C_B(qqq)v^3\int\delta^4(p-\sum_i p_i)\prod_{i=1}^{3}w(p_i)d^4p_i\ , \quad (5)$$

where $C(qqq)$ is the probability that three quarks chosen at random give a color singlet with the correct quantum numbers in a hadronic volume, v. The quark momentum distribution in the plasma is given as

$$w(p_i)d^4p_i=\frac{2gp_i}{(2\pi)^3}\delta(p_i^2-m_q^2)\{\exp\beta[\epsilon(p_i)+\mu]+1\}^{-1}\ , \quad (6)$$

where the statistical weight, $g=12$ for the u and d quarks, and $g=6$ for the strange quarks. $C(qqq)$ is computed using the spin-flavor wave functions for quarks in a baryon. The energy of the quark is $\epsilon(p_i)$ and satisfies the condition:

$$[\sum_i\epsilon(p_i)+BV]^2=(\sum_i p_i)^2+m_B^2\ . \quad (7)$$

Probability distribution functions for the mesons can be computed in a similar way. Assuming thermal equilibrium in the hadronic phase, the hadron momentum distribution is

$$W_N(p)\delta(p^2-m_N^2)d^4p=P(qqq)\frac{2g_NVp}{(2\pi)^3}\delta(p^2-m_N^2)\times(\exp\{\beta[\epsilon_N(p)+\mu_N]\}+1)^{-1}\ . \quad (8)$$

The energy of the hadron is ϵ_N, while the chemical potential of the hadron is μ_N. It is related to the quark chemical potential through the relation:

$$\mu_N=3\mu\ .$$

V denotes the plasma volume and g_N is the statistical weight factor for the baryon, N. Then for a nonrelativistic gas of hadrons, the mass spectrum is

$$\int W_N(p)\delta(p^2-m_N^2)d^4p=V\left[\frac{m_N}{2\pi\beta}\right]^{3/2}[\exp\beta(m_N+\mu_N)+1]^{-1}P(qqq)\ . \quad (9)$$

Equating this to the expression Eq. (5) derived earlier,

$$P(qqq)V\left[\frac{m_N}{2\pi\beta}\right]^{3/2}[\exp\beta(m_N+\mu_N)+1]^{-1}$$
$$=C(qqq)v^3\int\delta[(\sum_i p_i^2)-m_N^2]\prod_{i=1}^{3}w(p_i)d^4p_i\ , \qquad (10)$$

which determines $P(qqq)$.

The hadronization procedure described previously, although rigorous, has a serious computational drawback arising from the uncertainty in the hadronic radius, R. Since the hadronic volume goes as R^3, and therefore $P(qqq)$ varies as R^9, a small change in radius results in large variations in baryon and meson densities after hadronization. Furthermore, different hadrons can have different radii. This computational problem is resolved by demanding that the energy, baryon number, strangeness, and charge of the system be conserved during hadronization. Thus the baryon and meson densities in the rest frame of the plasma can be defined as

$$\rho_B=C_B\int\frac{1}{(2\pi)^3}d^3q\{\exp[\beta(\epsilon_B+\mu_B)]+1\}^{-1} \qquad (11)$$

and

$$\rho_M=C_M\int\frac{1}{(2\pi)^3}d^3q[\exp(\beta\epsilon_M)-1]^{-1}\ , \qquad (12)$$

respectively, where the meson energy $\epsilon_M=(q^2+m_M^2)^{1/2}$. The coefficients C_B and C_M are then adjusted to fit the energy, baryon number, strangeness, and charge densities at the transition point. While this simplified procedure is not as elegant as that described earlier, it is physically equivalent and much easier to compute. In the present calculation we include the N, Δ, Λ, and $\sum$ baryons and both the spin 0 and spin 1 octets of mesons. The final state interactions among the hadrons are not taken into account, but since the resonances, Δ, ρ, and others are explicitly considered, some part of this final state interaction is included.

III. RESULTS AND DISCUSSION

A. Antinucleon-nucleon annihilation

Experimental analysis[17] confirms that over 97 percent of the annihilation products of antinucleon-nucleon interactions at rest are pions and less than 3 percent are kaons. The pion and kaon momentum distributions[18] follow a Boltzmann form with characteristic temperatures, $T_\pi=128$ MeV and $T_K=84$ MeV indicative of the formation of a fireball. Statistical theory[19] predicts the total pion multiplicity in terms of a volume characterized by a fireball radius of 6.7 fm, which is unrealistically large compared to the Compton wavelength of a pion which is 1.4 fm. Pion interferometry[20] measurements confirm this interaction radius at somewhat higher energies to be 1.0 ± 0.3 fm. Bootstrapping with a single pionlike object[21] gives a multiplicity close to the measured value[22] of 5.1 ± 0.23 for an interaction radius of the pion Compton wavelength.

It is attractive to consider the s-wave annihilation process as one where all the constituents of the nucleon and antinucleon annihilate into a quark-gluon plasma of energy 1.87 GeV. The expansion and cooling of this plasma leads to hadronization. In this case of total annihilation the baryon number density and the chemical potential are both zero. The condition that the energy density $\epsilon=4B$ enables us to compute the volume of the plasma at the transition point. From our calculations this value is 2.8 fm^3 for $B^{1/4}=190$ MeV but increases to 7.0 fm^3 for the lower value of $B^{1/4}=150$ MeV. With the constraint that the sum of the meson energy densities should equal the total energy density, the number densities of the various mesons are determined from Eq. (12) and the calculation includes the pseudoscalar octet (π, K, and η), the pseudosinglet (η'), the pseudovector octet (ρ, K*, and ω), and the singlet (ϕ). As expected, the meson densities decrease as their masses increase, and apart from the pseudoscalar octet, only the ρ and ω densities are significantly high. All other higher mesons may be ignored without changing the results appreciably. Meson multiplicities are then obtained from the product of the plasma volume and the appropriate densities.

The results for the pion multiplicity $\langle n\pi\rangle$ are presented in Table I and show that the multiplicity does not depend sensitively on the coupling constant α_s nor on the saturation factor F. There is however a strong dependence on the bag constant, an increase of $B^{1/4}$ from 150 to 190 MeV reducing the multiplicity by about 15 percent. The calculated value is close to the experimental data which is encouraging in view of the fact that the application of statistical thermodynamics must be approximate when applied to such a small system. In such cases the limitation on phase space is important but difficult to compute in this format.

More than 25 to 40 percent of the pions are secondary, and of these, more than half are decay products of the ρ meson. The fraction of secondary pions increases with B through an increase in the plasma temperature. The variation of plasma temperature as a function of the different input parameters is shown in Fig. 1 for antiproton-proton annihilation. Extending our model to s-wave annihilation at higher energies we predict that the plasma volume must increase linearly with the total available energy, $\sqrt{s}$, since deconfinement is assumed to occur at a specified energy density. Qualitatively the meson multiplicity will therefore increase linearly with $\sqrt{s}$, while the momentum distribution remains unchanged. At higher energies of a few GeV the large relative momentum of the constituents may prevent equilibration. This lack of equilibration and the effects of peripheral annihilation will be reflected in a change in the rapidity distribution of the pions and a slower rate of increase of the multiplicity.

The neutral to charged ratio, $\langle n_{\pi^0}\rangle/\langle n_{\pi^+}\rangle$, for primary pions is unity, but the contribution from the decays of heavier mesons increases this ratio. For instance, η meson decay gives a neutral to a charged pion ratio of $\frac{71}{29}$. Detailed calculations finally show that $\langle n_{\pi^0}\rangle$ is about ten percent larger than $\langle n_{\pi^+}\rangle$.

The ratio $\langle n_k\rangle/\langle n_\pi\rangle$ depends almost linearly on the saturation factor F, and the values of $F=0.2$ with the $B^{1/4}$ value of 150 MeV and $F=0.1$ with $B^{1/4}=190$ MeV fit the data. These values are consistent with the esti-

TABLE I. Results for the $n\bar{n}$ annihilation through the formation of a quark-gluon plasma.

α_s	F	T (MeV)	$\langle n_\pi \rangle$ primary	$\langle n_\pi \rangle$ secondary	$\langle n_\pi \rangle$ total	$\langle K/\pi \rangle$	$\langle \pi^0/\pi^+ \rangle$
$B^{1/4}=150$ MeV							
0.5	0.1	122.9	3.42	1.06	4.48	0.017	1.09
0.5	0.2	121.9	3.38	1.05	4.43	0.028	1.08
0.6	0.1	127.5	3.15	1.21	4.35	0.020	1.10
0.6	0.2	128.9	3.13	1.17	4.30	0.032	1.10
$B^{1/4}=190$ MeV							
0.5	0.05	156.4	2.32	1.69	4.01	0.026	1.13
0.5	0.1	155.7	2.29	1.67	3.96	0.035	1.12
0.6	0.05	164.2	2.13	1.80	3.93	0.030	1.14
0.6	0.1	163.2	2.12	1.76	3.88	0.039	1.13

mates of Rafelski and Mueller[13] for a plasma lifetime of 10^{-23} s. The momentum distribution of the pions, shown in Fig. 2, is in rather good agreement with experiment provided that the bag constant is 150 MeV. Taking into account the finite phase space would reduce the higher momentum components of the momentum distribution to the experimental values with a cutoff at about 0.96 GeV/c.

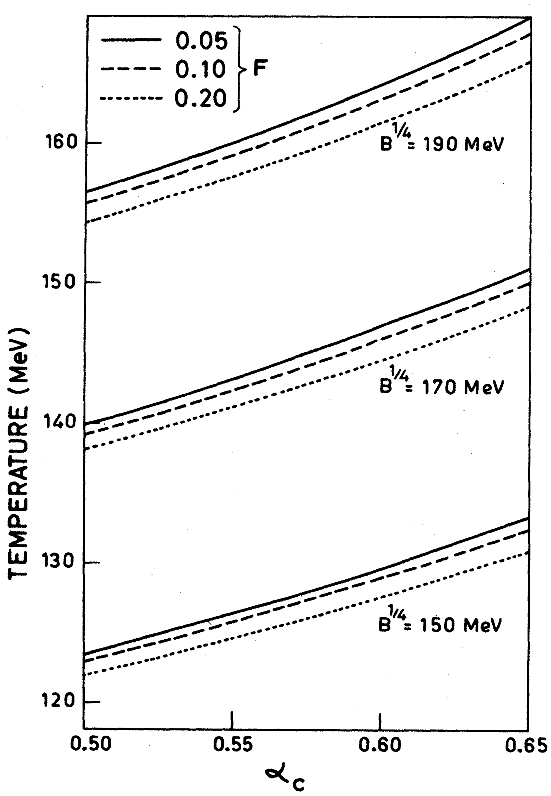

FIG. 1. Plasma temperature for various values of the bag parameter B, the saturation factor F, and the QCD coupling constant α_s, in antiproton-proton annihilation.

B. Central heavy ion collisions

Ultrarelativistic heavy ion collisions have been the favorite hunting ground[23] for searches for the quark-gluon phase, and the possible signals for the formation of such a phase have been discussed at length. It is likely that even at lower energies a quark-gluon plasma is formed and the production of strange particles below the nn threshold with high transverse momentum could be a characteristic signal of plasma formation.

The center of mass energy for the nucleon-nucleon collision in the central collision of two equal mass nuclei is

$$E_{\mathrm{c.m.}} = 2m(1+T_{\mathrm{lab}}/2m)^{1/2}\ , \tag{13}$$

where T_{lab} is the incident energy per nucleon in the laboratory. The energy of each nucleon being $E_{\mathrm{c.m.}}/2$, the energy density in the central region of density ρ is

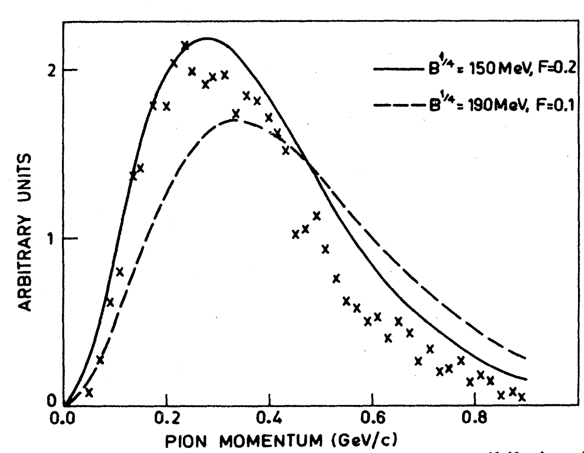

FIG. 2. Pion momentum distribution in $n\bar{n}$ annihilation for two values of the bag parameter and the saturation factor. The data are from Ref. 18.

TABLE II. Nucleon density n, plasma temperature T, and compression factor CF for heavy ion collisions at various energies with $\alpha_s = 0.6$ and $F = 0.2$.

Energy (GeV/nucleon)	$B^{1/4}=150$ MeV n (fm^{-3})	T (MeV)	CF	$B^{1/4}=190$ MeV n (fm^{-3})	T (MeV)	CF
1.5	0.104	110.0	1.23	0.269	121.5	3.17
1.8	0.100	112.0	1.18	0.258	125.5	3.04
2.1	0.096	113.5	1.14	0.248	129.0	2.92
2.4	0.092	114.5	1.09	0.238	132.0	2.82
2.7	0.089	115.5	1.06	0.231	134.5	2.72

$$\epsilon = m(1+T_{\text{lab}}/2m)^{1/2}\rho \ . \tag{14}$$

Monte Carlo calculations[24] on the central collision of heavy ions of mass 40 at an incident energy of 2 GeV/nucleon show that soon after the interpenetration begins (10^{-23} s), a high density of 3—4 ρ_0 builds up in the central region and that kinetic energy of the participating nucleons is to a large extent equilibrated. The consequent energy density is around 1 GeV/fm^3, not much lower than that within a hadron. The transition to the quark-gluon phase is just possible.

To avoid counting the u and d quarks separately we consider identical colliding nuclei of $A = 2Z$. The energy density, temperature, and chemical potentials for these quarks at the transition point are computed as indicated earlier. Our coalescence model then determines the density of the π, K, and ρ mesons, the nucleons, and the Δ, Λ, and Σ baryons. Higher mass hadrons are neglected as they do not alter the results significantly.

The dependence of the nuclear density and plasma temperature is given in Table II along with values of a compression factor which is defined as the minimum compression of hadronic matter required to form the plasma. The computed quantities are not sensitive to the incident energy nor to the values of α_s and F. The dependence on $B^{1/4}$ is however interesting. At 150 MeV practically no compression is necessary to form the plasma phase and all the nuclear matter participates in plasma formation. But if a value of $B^{1/4}=190$ MeV is assumed, a compression of 3 is required for the transition. Both cascade[24] and hydrodynamic[25] calculations indicate that compressions of at least this magnitude occur in heavy ion collisions of 2 GeV/nucleon energy. As the bag constant increases to 190 MeV from 150 MeV, the plasma temperature increases from 105 to 120 MeV. It is comforting to note that the temperatures extracted[26] from measured proton, kaon, and pion spectra in central collisions at this energy are in this range.

It is however difficult to estimate the volume of the plasma, and therefore, the absolute cross sections and the multiplicities cannot be estimated. We therefore compute only particle production ratios for pion, kaon, and lambda to the proton cross section at various incident energies, and the results are shown in Fig. 3. The strange hadron ratios are generally larger for the smaller values of $B^{1/4}$ and increase appreciably with energy. For example, the K^-/p ratio increases by a factor of 3 ($B^{1/4}=150$ MeV) or a factor of 8 ($B^{1/4}=190$ MeV) as the energy increases from 1.5 to 3.0 GeV/nucleon. Recently several experiments[26—28] which have measured Λ, K^+, and K^- differential production cross sections in heavy ion reactions indicate that the Λ and K^+ particles are isotropic in the center-of-mass frame and their spectra correspond to a fireball of temperature 120 MeV (Fig. 4). This is in agreement with our computations on the quark-gluon coagulation model. The experimental multiplicity ratios are $\Lambda/\pi^- = 6.7\times10^{-3}$, $\text{K}^-/p = 4\times10^{-3}$, and $\text{K}^+/\text{K}^- = 50\times10^{-3}$ which are close to our results with $B^{1/4}=150$ MeV. An estimate of Λ-production cross section for the Ar + KCl collision at 1.8 GeV/nucleon can be obtained by assuming a plasma volume as given by the maximum impact parameter of 2.4 fm determined[26] by the central trigger of the experiment. This volume corresponds to about 20 nucleons participating in plasma formation and the computed value of 6.1 mb is then in agreement with the experimental result of 7.6 ± 2.2 mb. For K^+ production in the Ne + NaF reaction at 2.1 GeV/nucleon there was no selection of central collisions. Choosing a reasonable maximum impact parameter of 3 fm, which would be the radius of a nucleus with $A = 20$, the plasma volume of 35 fm^3 gives a calculated cross section of 16 mb, in satisfactory agreement with the measured 23 ± 8 mb.

IV. CONCLUSIONS

The preceding discussion indicates that the results of our calculations on the coagulation model are in fairly good agreement with data. However estimates[29] of strange particle production exist where only a hadronic fireball is considered and where the existence of a quark-gluon phase is denied,[30] but these estimates require as input, free particle cross sections. Off-shell production amplitudes must necessarily be neglected in such a formulation and these would be important in the time scales of these reactions at the energies considered here. The off-shell contribution is to some extent included when the constituents of the nucleons are explicitly considered.

The similarity between the antinucleon-nucleon annihilation process and the central collision of heavy ions also makes the quark-gluon coagulation model attractive, particularly as the bag pressure, QCD coupling constant, and the saturation factor are almost identical and reasonable in both the reactions considered.

At the moment there are no definite experimental signatures that identify the formation of a quark-gluon phase rather than a hadronic fireball, but the energy dependence of strange particle production may provide some indication. In our model, the $S=-2$ (Ξ) baryon density is between 0.4×10^{-6} and $1.6\times10^{-6}/\mathrm{fm}^3$ which leads to a production cross section of 10—40 μb in the Ar + KCl collision at 1.8 GeV/nucleon. Antibaryon production has also been proposed as a signal of plasma formation but most of the antibaryons even if formed will be annihilated with a cross section of 100 μb within the fireball and will not be detected by the experiment.

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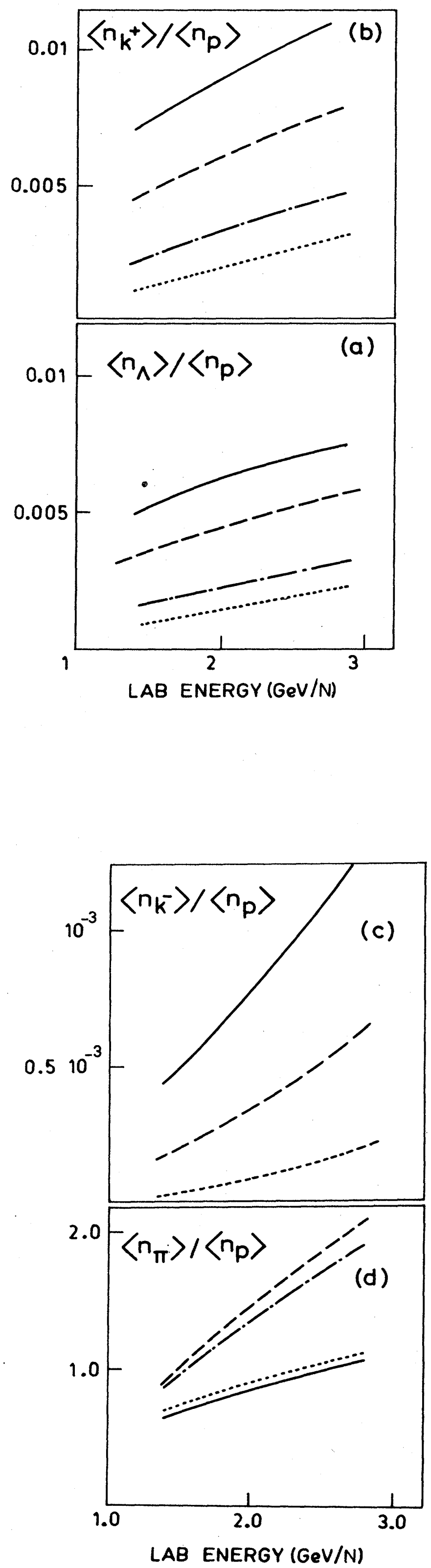

FIG. 3. Ratios of pion, kaon, and lambda to proton densities as a function of incident energy. (i) $B^{1/4}=190$ MeV, —— $F=0.2$; — — — $F=0.1$. (ii) $B^{1/4}=150$ MeV, —·—·— $F=0.2$; ··· $F=0.1$. For all curves, $\alpha_s=0.5$.

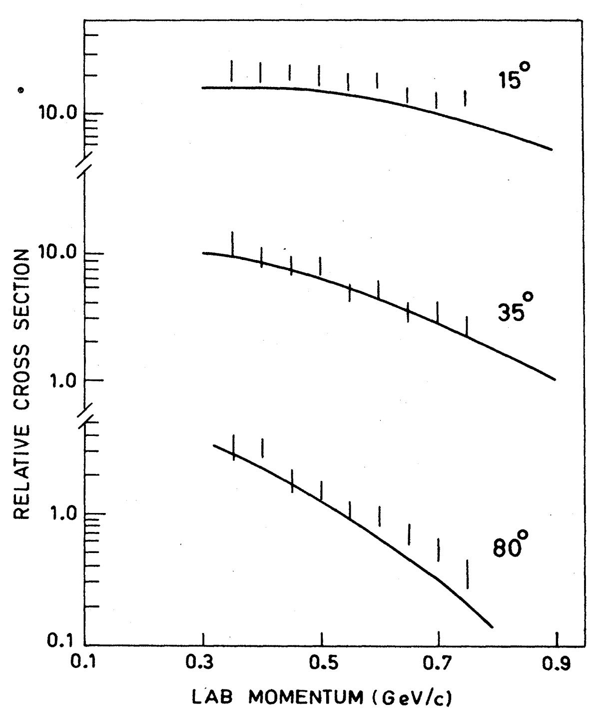

FIG. 4. K^+ spectra at three angles for $B^{1/4}=150$ MeV, $\alpha_s=0.6$, and $F=0.2$ compared to experimental data of Ref. 28.

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