Possible production of exotic baryonia
in relativistic heavy–ion collisions

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Abstract

Properties of a hypothetical baryonium with the quark content $(udsuds)$ are discussed. The MIT bag model predicts its mass to be unexpectedly low, approximately 1210 MeV. Possible hadronic decay modes of this state are analyzed. Ultrarelativistic heavy–ion collisions provide favorable conditions for the formation of such particles from the baryon–free quark–gluon plasma. We estimate multiplicities of such exotic baryonia on the basis of a simple thermal model.

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Recent experimental observations of exotic baryons, namely the $\Theta^+(1540)$ and $\Xi^{--}(1862)$ pentaquarks, are of great interest for understanding strong interactions. On the other hand, many predictions of exotic hadrons, made within different effective models of QCD are still waiting for their experimental verification. For instance, the dibaryon state with the quark content $(u^2d^2s^2)$ has been predicted using the MIT bag model and not found yet. In Refs. this model was used to calculate masses of exotic multi-quark-antiquark states, in particular, the pentaquarks $(q^4\bar{q})$ and heptaquarks $(q^5\bar{q}^2)$. Later, low-lying pentaquark states have been predicted within the chiral soliton model.

In Ref. binding energies of multi-$q\bar{q}$ systems with different flavor compositions have been studied within a completely different approach. Namely, the equation of state of $q\bar{q}$ matter at fixed relative concentrations of antiquarks was calculated within the generalized Nambu–Jona-Lasinio model. It was shown that such systems have bound states at zero pressure. The maximum binding energy per particle was obtained for systems with equal number of quarks and antiquarks, so-called ”mesoballs”. In these systems the repulsive vector interaction vanishes (at least on the mean–field level) and the binding energy is determined by a competition between the scalar attraction and the kinetic energy of quarks and antiquarks. The kinetic (Fermi) energy is lowered by transforming a fraction of light $q\bar{q}$ pairs into strange $s\bar{s}$ pairs. Such flavour symmetrization is also favored by the instanton–induced flavor-mixing interaction. As a consequence, the binding energy per quark–antiquark pair is predicted to be maximal for systems with large hidden strangeness. In these mean–field calculations the flavor and color correlations are considered only on average, and therefore, clusterization of $q\bar{q}$ matter is ignored. Possible formation of multi-quark-antiquark systems was suggested also in Ref. in connection with deeply bound antibaryon states in nuclei. The role of attractive color configurations of quarks and gluons and possible existence of their bound states in the quark–gluon plasma have been discussed in Refs.

Recently one of us reconsidered the systematics of multi-$q\bar{q}$ clusters by using the MIT bag model with the same set of parameters as in Ref. In particular, masses of exotic systems with the quark content $(u\bar{u})^m(d\bar{d})^n(s\bar{s})^k$, where $m + n + k = 3, 4, 5, 6,$
have been calculated. We use the name exotic baryonia for states with \( m + n + k = 3 \) in order to distinguish them from the \( BB \) bound states discussed long ago within the mesonic picture of strong interaction (see Refs. \[13, 14\]). Recent measurements \[15\] of the \( p\bar{p} \) mass spectrum in \( J/\psi \to \gamma p\bar{p} \) decays seem to give experimental evidence in favor of the subthreshold \( NN \) bound state with mass \( m \simeq 1859 \text{ MeV} \). This observation has renewed interest in the issue of baryonia. For instance, in Ref. \[16\] the \( \Lambda\bar{\Lambda} \) baryonium with mass \( m \simeq 2200 \text{ MeV} \) has been proposed.

In our opinion, the above mentioned loosely bound states are not the lowest energy states of baryonia. According to the MIT bag calculations of Refs. \[12, 17\] the lightest \( q^3\bar{q}^3 \) state is a flavor symmetric pseudoscalar \( (J^P = 0^-) \), with the quark content \( (u\bar{u}d\bar{d}s\bar{s}) \). Such state has zero isospin, zero strangeness and positive G-parity. For brevity we call it \( Y \)-particle. Its mass is predicted to be unexpectedly small, \( m_Y \simeq 1214 \text{ MeV} \), i.e. significantly below the \( \Lambda\bar{\Lambda} \) threshold \( 2230 \text{ MeV} \). It has such a low mass because of the unique arrangement of quarks and antiquarks leading to a very strong color-magnetic interaction \[28\]. Within the Jaffe approximation \[5\], the mass-dependent factors in the interaction terms are replaced by average values. Then the color-spin interaction reduces to terms that can be expressed via Casimir operators of \( SU(6)_{cs} \), \( SU(3)_c \), and \( SU(2)_s \). To obtain antisymmetric wave functions, the flavor \( SU(3)_f \) symmetry representation should be conjugate to the \( SU(6)_{cs} \) representation. The larger the symmetry of the \( 3q \) or \( 3\bar{q} \) wave functions in \( SU(6)_{cs} \), the greater is the color-magnetic attraction. The three-quark states that are completely symmetric (i.e., the \([3]\)-states) do not contain \( SU(3)_c \) singlets \[7\] and hence, they can not exist as free particles. Only by coupling to three antiquarks with the same symmetry, one can produce a color singlet state. Thus, the \((uds)\) and \((\bar{u}\bar{d}\bar{s})\) combinations in the \( Y \)-particle are not color singlets, unlike in the quasimolecular \( \Lambda\bar{\Lambda} \) state.

Let us discuss possible decay modes of the \( Y \)-particle. First of all we note that such simple channels as \( Y \to K\bar{K} \) and \( Y \to n\pi \) with an even number of secondary pions are forbidden due to the parity and angular momentum conservation. The decays into odd numbers of pions should be suppressed because of the G-parity violation. Thus, we expect that only the 3-body decay channels \( Y \to \pi K\bar{K} \) and \( Y \to 2\pi\eta \) may contribute
significantly to the width. These final states can be obtained by simple rearrangement of quarks from the initial state. However, it is difficult to calculate the corresponding matrix elements. If they are of the same order, the relative contribution of these two channels is controlled by the phase space available for final particles. The relativistic phase space volume corresponding to the decay of a particle with mass $m$ into $n$ particles with masses $m_1, \ldots, m_n$ is defined as

$$R_n = \int \prod_{i=1}^{n} \frac{d^3 k_i}{k_i^0} \delta^4(P - \sum_{l=1}^{n} k_l), \quad (1)$$

where $P^\mu = (m, 0)^\mu$ and $k_i^\mu = (\sqrt{m_i^2 + k_i^2}, k_i)$ are, respectively, the 4-momenta of initial and final particles (in the rest frame of the decaying particle). The direct calculation gives

$$\frac{R_3(Y \to \pi K \bar{K})}{R_3(Y \to 2\pi \eta)} \approx 0.0543. \quad (2)$$

One can see that the $Y \to \pi K \bar{K}$ decay channel is strongly suppressed due to closeness to the threshold. On the other hand, the $\eta$ meson has smaller hidden strangeness than the $K \bar{K}$ pair, and thus, the $Y \to 2\pi \eta$ decay mode may be additionally suppressed by the OZI rule. From the experimental viewpoint, using this decay channel is problematic because of the small width of $\eta$ meson. We would like also to point out that, similarly to $\eta$ meson, the $Y$-particle should have a significant $\gamma \gamma$ width, most likely, in the keV range. Therefore, the observation of exotic $Y$-baryonium in $pp$ or $\pi p$ reactions should be possible by measuring the $\pi K \bar{K}, 2\pi \eta$ and $\gamma \gamma$ invariant mass spectra. Because of the pseudoscalar nature of the $Y$-particle we do not expect its formation in $e^+e^-$ collisions.

It is worth noting that in the mass region of interest there exists one meson, the pseudoscalar-isoscalar $\eta(1295)$, which cannot be interpreted as a conventional $q\bar{q}$ state. It was seen in the reactions $\pi^- p \to \eta \pi^+ \pi^- n (\eta \to \gamma \gamma)$ and $\pi^- p \to K^+ K^- \pi^0 n$ studied in Refs. \cite{19} and \cite{20}, respectively. The width of this meson is estimated to $50 \div 70$ MeV, but its decay channels are still poorly known. In principle, $\eta(1295)$ can be identified with the $Y$-particle. However, such interpretation would require a more detailed analysis of the $\eta(1295)$ and $Y$ decay modes.

We believe that $Y$-baryonium can be naturally produced in ultrarelativistic heavy–ion collisions e.g. at the RHIC in Brookhaven. It is expected that a nearly baryon–free
quark–gluon plasma is formed at an intermediate stage of such collisions. Then multi–quark-antiquark clusters should be also formed with a certain probability. Analysis of experimental data shows that ratios of hadron multiplicities observed in central collisions of nuclei in a broad range of bombarding energies are well reproduced within a simple thermal model. According to this model the multiplicity  $N_i$ of hadron species $i$ with mass $m_i$ and spin $J_i$ is expressed as

$$N_i = \frac{(2J_i + 1)V}{(2\pi)^3} \gamma_s^{n_s} \int d^3p \left[ \exp \left( \frac{\sqrt{m_i^2 + p^2 - \mu_i}}{T} \right) \pm 1 \right]^{-1},$$

where $+(-)$ corresponds to fermions (bosons). It is assumed that all hadrons are formed in an equilibrated system characterized by volume $V$ and temperature $T$. In Eq. (3) $\mu_i$ stands for the chemical potential of corresponding hadrons. Following Refs. [21, 23] we take into account possible deviations from chemical equilibrium for hadrons containing nonzero number ($n_s$) of strange quarks and antiquarks, by introducing a strangeness suppression factor $\gamma_s$. On the other hand, approximately the same degree of agreement with experimental data is achieved in Ref. [22] assuming $\gamma_s = 1$. Numerical values of the parameters used in Refs. [21] (set B) and [22] for central Au+Au and Pb+Pb collisions at different c.m. energies $\sqrt{s}$ are given in Table I.

Below we estimate multiplicity of $Y$-particles, which might be produced in central A+A collisions, proceeding from the ratio $N_Y/N_\phi$ predicted by the thermal model. This choice is motivated by two reasons. First, $Y$ and $\phi$ contain the same number of $s\bar{s}$ pairs ($n_s = 2$)

### Table I: Fitted parameters of thermal model for central Au+Au and Pb+Pb collisions at AGS, SPS and RHIC energies.

| reaction | $\sqrt{s}$ (AGeV) | $T$ (MeV) [21] | $\mu_B$ (MeV) [21] | $\gamma_s$ [21] | $T$ (MeV) [22] | $\mu_B$ (GeV) [22] |
|----------|-------------------|----------------|---------------------|-----------------|----------------|-------------------|
| Au+Au    | 4.88              | 119.1          | 578                 | 0.763           | 125            | 540               |
| Pb+Pb    | 8.87              | 145.5          | 375.4               | 0.807           | 148            | 400               |
| Pb+Pb    | 12.4              | 151.9          | 288.9               | 0.766           |                |                   |
| Pb+Pb    | 17.3              | 154.8          | 244.5               | 0.938           | 170            | 255               |
| Au+Au    | 130               | 165            | 41                  | 1.0             | 176            | 41                |
and, therefore, $\gamma_s$ factors cancel in the ratio $N_Y/N_\phi$. Second, their masses are close and as a consequence, the canonical suppression factors [22], which become important at low energies, are also approximately cancelled.

In the Boltzman limit, taking into account that $\mu_Y = \mu_\phi = 0$ and substituting $J_Y = 0, J_\phi = 1$ into Eq. [3], one has:

$$\frac{N_Y}{N_\phi} = \frac{1}{3} \left( \frac{m_Y}{m_\phi} \right)^2 \frac{K_2(m_Y/T)}{K_2(m_\phi/T)},$$

(4)

where $K_2(z)$ is a modified Bessel function of the second order.

Using the parameters of Table I and multiplicities of $\phi$ mesons given in Refs. [21, 22, 26], we calculated absolute multiplicities of $Y$-particles in central Au+Au (AGS, RHIC) and Pb+Pb (SPS) collisions. The results are shown in Fig. 1. The discrepancy between two

![Graph showing average multiplicities of $Y$-particles in central Au+Au and Pb+Pb collisions at different c.m. bombarding energies. The solid and dashed curves correspond to two sets of parameters suggested, respectively, in Refs. [21] and [22].](image-url)
estimates at $\sqrt{s} = 17.3$ GeV is explained by two different data sets on $\phi$ multiplicities reported by the NA49 \cite{24} and NA50 \cite{25} collaborations. One can see that at energies $\sqrt{s} \gtrsim 10$ AGeV the predicted $Y$-multiplicities are rather high, $N_Y \gtrsim 1$, which makes their experimental observation feasible. As discussed above, the decay modes, $Y \to \pi K\bar{K}$, $Y \to 2\pi\eta$ and $Y \to 2\gamma$ could be used for experimental identification of these particles.

Recently in Ref. \cite{27} a similar consideration was used to estimate the abundance of $\theta^+$ pentaquarks in nuclear collisions at RHIC. About one $\theta^+$ per unit rapidity was predicted per central Au+Au collision. According to our estimates, due to the lower mass of $Y$-particles, their yield at midrapidity should be several times larger.

Finally we want to emphasize that the MIT bag model is a phenomenological model which nevertheless contains many essential features required for a description of hadrons. Its accuracy for systems with large number of quarks and antiquarks is unknown. However, the large attractive color–magnetic interaction suggests that the $Y$-baryonium state should exist, even if the calculated mass might be only approximate.

In conclusion, if a strongly bound baryonium state predicted by the MIT bag model exists and its decay width is not too large, it can be produced at the hadronization of a quark–gluon plasma. Using a simple thermal model we have estimated yields of these particles in central collisions of relativistic heavy ions. Expected multiplicities are large enough, $\sim 5$ per event at RHIC, for their detection. This observation would not only prove the existence of a new exotic hadron, sextaquark, but also provide a very strong evidence in favor of the quark–gluon plasma formation.

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