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DO NUCLEONS DISSOLVE IN GIANT NUCLEI?

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Abstract. An interesting and until now unanswered question is how quarks, confined within a bag, interact when their number is very large. Can their residual interaction be understood in terms of perturbation theory or is it basically non-perturbative? As remote as any other possibility is to assume the quarks to be free within the bag, when the total baryon number corresponds to that of giant nuclei. We investigate here such giant quark bags with baryon numbers up to 1000. Experimental consequences and some speculations in connection with the recently observed anomalous positron production are discussed.

I. INTRODUCTION

On the nowadays as fundamental accepted level of QCD the nucleon-nucleon interaction appears to be an effective description of forces acting between the truely elementary particles, the quarks. The hadrons, and among them also the nucleons, appear to be composite objects containing small numbers of quarks trapped within a limited region of space. Due to the complex algebraic structure of the underlying symmetry group $SU(3)_c$, which, for example, has the consequence that the eight mediators of this interaction, the gluons, interact with each other, even rather elementary calculations turn out to be possible only by help of large computers. Model builders have, however, anticipated the results of future exact calculations, and developed the so called "bag" picture of hadrons[1]. In this approach the strong interaction is thought to be split into two parts. The first one, which describes the non-perturbative and non-abelion part responsible for the vacuum structure and the confinement of quarks, cannot be calculated explicitly and is built in "by hand". Excitations of the ground state are bubble-like inhomogenities in the vacuum condensate, filled with quarks which can be also regarded as local minima in a kind of self-consistent potential analogous to the effective pairing potential known from the Landau-Ginsburg theory of superconducting materials[2]. Hadrons are nothing else than such holes in the vacuum-medium digged by the quarks themselves. First order perturbation theory (one-gluon exchange) turns out to be sufficient for the residual part of the interaction, when baryons and, to some degree also mesons, with a small number of quarks are considered.

Models based on this semi-phenomenological point of view exist in various degrees of sophistication (for example the MIT-bag[3], the soliton-bag[4], the chiral bag[5], the hybrid bag[6], etc.[1]). In all of them the confining effect of the surrounding vacuum is achieved through a substantial increase of the quark mass across the boundary of the bubble.

Conventionally, atomic nuclei are considered to be composed of nucleons, treated as its fundamental (pointlike) components which do not change

their properties when embedded in nuclear matter and interact via some quite involved and empirically determined interaction. This point of view has lead to substantial success in explaining low energy data and the static properties of atomic nuclei. Recent experiments revealing the high momentum (i.e. small distance) structure of nuclear matter[7] force us, however, to modify this point of view, since the internal structure of nucleons embedded in nuclei seems to be altered. The degree of modification of the form factor even increases with the size of the surrounding nucleus.

In the framework of the bag-models the above-mentioned effect on the nucleon structure can be well understood[8]. Although the quarks feel "uncomfortably" outside the bag, when the exterior is a simply homogeneous condensate, i.e. when the hadron (nucleon) is isolated, they may leak out, if the vacuum becomes structured in its vicinity due to the presence of other hadrons. Moreover, since the condensate itself responds to the quark structure in a self-consistent way, the inhomogenities themselves should change, too. The probability for a coloured quark to be outside the free nucleon thus increases, i.e., the confinement radius becomes effectively larger, the nucleons overlap and colour-conductivity sets on[9]. With increasing delocalisation of the quarks a colour band structure can develop, resembling band effects known from the elementary physics of conducting crystals. Corresponding calculations for periodic bag structures have been done in the framework of the hybrid bag model[10] and of the MIT-bag model[11].

In this paper we do not wish to add another piece of work to the already existing vast literature on the explanation of the EMC-effect[8], but take the picture discussed above seriously and extrapolate it to very large nuclei: if the colour-conductivity within nuclei increases with their increasing size, and if this increase is connected with a continuos depletion of the inhomogenities in the "condensate field" (i.e. the σ-field in the language of the soliton-bag model), then in giant nuclei with nucleon number A of the order of several hundreds this field should vanish completely in the interiour, turning the giant nucleus into a cold quark-gas. These objects, which we call giant quark nuclei (GQN), are

investigated here[12]. The question arises, of course, whether the perturbative treatment of the residual quark-quark interaction is applicable also when the number of quarks and the size of the "hadron" are very large. A possibility as remote as any other is to assume that the importance of the residual interaction decreases in this case. One can argue against this point of view, but lacking any conclusive prescription how to handle the residual interaction, a calculation based on our assumption can at least serve as a basis for future "better" calculations.

The transition of ordinary nuclear matter into a colour conductor, which we have in mind here, seems to be rather analogous to similar phenomena observed in solid state physics, when clusters of atoms are studied with respect to the development of collective properties like conductivity or superfluidity. There it is found that such properties do not appear, if the number of participating atoms is relatively small (10^2, 10^3, 10^4) but that they set on beyond a certain cluster size[13].

There is still another motivation for our investigation on Giant Quark Nuclei, emerging from the very exciting studies of the decay of the electron-positron vacuum in supercritical electric fields generated by two very heavy ions in collisions just above the Coulomb barrier[14]. Many experimental hints point towards the discovery of positrons spontaneously produced in this fundamental process and also towards the existence of giant nuclei. The intimate link between these two concepts is due to the appearance of a main and several secondary positron resonance lines at various energies and in different "windows", which possibly indicate the existence of surprisingly long-living nuclear objects. Even though phenomenological studies of unified nuclear systems show that nuclear molecular states are not unlikely[15], giant nuclei seem to be difficult to justify on the basis of conventional nuclear physics[16] or even in Meson Field Theory[17]. Nevertheless, the strange scaling property of the experimentally observed and so far principally investigated positron peaks definitely requires a new type of nuclear structure.

The paper is organized in the following way: in Section II we introduce a simple model to estimate the binding energy per nucleon of Giant Quark

Nuclei in the framework of the MIT-bag approach, and compare the results with those obtained from the extrapolation of the empirical Bethe-Weizsäcker formula fitted to known nuclei. In Section III we discuss the physical content of this calculation and its consequences in the context of the bulk of known data. Finally, speculations about experimental consequences seem appropriate.

II. MODEL FOR GIANT QUARK NUCLEI

We consider now a nucleus with the mass number A, in which quarks are not clustered into nucleons, but move freely throughout the interiour of the whole nucleus. The confining force keeping the quarks from leaving the nucleus can be understood as the action of the surrounding "true" QCD-vacuum in which coloured objects may not propagate unless they are combined into colour singlets. Such nuclei (or hadrons in general) form "bubbles" in the "true vacuum" filled with quarks. These bubbles correspond to local collective vacuum-excitations of the collective condensate field. The amount of energy, which has to be invested to create such an excitation, is proportional to the volume V of the bubble. The corresponding energy density B is a universal constant, which has to be fitted to experiments, as long as we are not able to connect it with the QCD scale parameter by exact calculations (see, however, Ref. 18). In the MIT-approach[3] such bubbles are described as cavities with sharp boundaries, in which the quarks (and eventually also gluons) move freely obeying the free equations of motion (we neglect here, as discussed above, the residual interaction, i.e. the one-gluon exchange):

$$i\gamma_\mu \partial^\mu \psi(x) = 0, \tag{1a}$$

where ψ is a four-component spinor, cf. eq. (30).

On the edges, however, their wave functions have to satisfy the boundary conditions (n^μ is the surface normal vector)

$$i\gamma_\mu n^\mu \psi = \psi \qquad \text{on the surface} \tag{1b}$$

$$-\tfrac{1}{2} n^\mu \partial_\mu \bar{\psi}\psi = B \qquad \text{on the surface} \tag{1c}$$

which mimic the infinite discontinuity of the quark mass across the boundary. This leads to a discrete single particle energy (and momentum) spectrum, displayed in Table I for massless quarks in the convenient dimensionless form

$$\omega_{\kappa n} = RE_{\kappa n}, \tag{1d}$$

where $E_{\kappa n}$ is the eigenfrequency of the quark level with the Dirac quantum number κ and the principal quantum number n, and R is the radius of the spherical cavity.

Adding a nucleon into the bag corresponds to the addition of two up- and one down-quark for a proton, and of one up- and two down-quarks for a neutron. Whereas the colour interaction responsible for the confinement does not depend on flavour, the electromagnetic charge becomes important when considering the effects of the Coulomb energy. This will be done schematically using the following approximation. Suppose the nucleus be a homogeneously charged sphere of radius R. The Coulomb energy of such a sphere (total charge Ze) is

$$E_c' = 3\alpha Z^2/5R, \tag{2}$$

where $\alpha = e^2/4\pi = 1/137$ is the fine-structure constant. Now the up- and down-quarks (their number be Z_u and Z_d, respectively) carry the fractional electric charges 2e/3 and -e/3, such that the required total charge of our nucleus is

$$Z = 2/3\ Z_u - 1/3\ Z_d \tag{3}$$

and the number of neutrons A-Z. Because

$$Z_u + Z_d = 3A, \tag{4}$$

$\omega_{\kappa n}$	κ		$\omega_{\kappa n}$	κ	
2.04278694273	-1	6	12.33348062758	5	822
3.20391876726	-2	18	12.70944981993	-4	846
3.81153864777	1	24	12.90161883497	-11	912
4.32730291250	-3	42	12.96372863983	-7	954
5.12311062410	2	54	13.00886854595	3	972
5.39601611785	-1	60	13.19690694178	-2	984
5.42952237092	-4	84	13.31559357684	1	990
6.37113691533	3	102	13.35096234860	9	1044
6.51788754767	-5	132	13.58486809101	6	1080
6.75780501185	-2	144	13.95153160392	-12	1152
7.00203329571	1	150	14.01249543641	-5	1182
7.58129857678	4	174	14.14515954635	-8	1230
7.59634581701	-6	210	14.36722103420	4	1254
8.05957827856	-3	228	14.47248193811	10	1314
8.40758707164	2	240	14.60960858443	-3	1332
8.57755878461	-1	246	14.78940655650	2	1344
8.66730055378	-7	288	14.81738110856	7	1386
8.76571372299	5	318	14.88782748573	-1	1392
9.32191544672	-4	342	14.99890082782	-13	1470
9.73233713563	-8	390	15.29173705254	-6	1506
9.75354968367	3	408	15.31495554700	-9	1560
9.93122588787	6	444	15.58709297961	11	1626
10.00419347084	-2	456	15.69559853618	5	1656
10.16332073509	1	462	15.98579945042	-4	1680
10.55591883243	-5	492	16.03430698797	8	1728
10.79256324954	-9	546	16.04402993356	-14	1812
11.05862603181	4	570	16.21798518269	3	1830
11.08213825437	7	612	16.36909138309	-2	1842
11.37644149760	-3	630	16.46389560006	1	1848
11.61203022487	2	642	16.47485855825	-10	1908
11.73650395934	-1	648	16.55150457940	-7	1950
11.76836247041	-6	684	16.69571662752	12	2022
11.84878712545	-10	744	16.99980620134	6	2058
12.22135972278	8	792	17.08716811425	-15	2148

Table I: Cavity modes $\omega_{\kappa n} \leq 17$ in ascending order. The Dirac quantum numbers κ are also displayed as well as the number of quarks with a particular flavour, filling all levels up to the given one.

$$Z_u = A + Z \tag{5a}$$

$$Z_d = 2A - Z. \tag{5b}$$

After inserting (3) into (2) we obtain the Coulomb energy

$$E_c' = 3\alpha/5R\ (2Z_u/3 - Z_d/3)^2 \tag{6}$$

of a homogeneously charged sphere consisting of Z_u up- and Z_d down-quarks. In this expression, however, also the Coulombic self-interaction of each quark is counted. This can be renormalized away by replacing

$$Z_i^2 \rightarrow Z_i(Z_i - 1) \qquad i = u,d \tag{7}$$

in eq. (6), thus yielding

$$E_c = \alpha/15R\ [4Z_u(Z_u-1) + Z_d(Z_d-1) - 4Z_uZ_d] \equiv \omega_c/R. \tag{8}$$

Insertion of the Coulomb energy makes the proton heavier then the neutron. We can cure it by realizing that the mass of the down-quark is larger than that of the up-quark. A reasonable assumption, which is consistent with standard bag calculations, is $m_u = 0$, $m_d = 5$ MeV. This small mass leads to a modification of the boundary condition for the quark wave-functions[3]. If we expand the corresponding equation around $\mu = 0$, in first order of $\mu = m_dR$ the arising correction of the momentum eigenvalue is estimated to be

$$\omega_{\kappa n}(\mu) = \omega_{\kappa n}(0) + \tfrac{1}{2}\mu/[\omega_{\kappa n}(0)+\kappa] \tag{9a}$$

and for the energy we obtain as usually

$$\varepsilon_{\kappa n}(\mu) = [\omega_{\kappa n}{}^2(\mu) + \mu^2]^{\frac{1}{2}}. \tag{9b}$$

With these considerations the total energy of the bag is simply[3]

$$E_{tot}(A,Z,R) = \Sigma\, \varepsilon_{\kappa n}/R + E_c(Z_u(A,Z),Z_d(A,Z)) + 4\pi BR^3/3. \qquad (10)$$

The sum in the first term corresponds to the kinetic energy of the quarks and runs over all occupied modes, the second term is the renormalised Coulomb energy (8), and the last one the condensation or volume energy with the pressure taken to be[19] $B = (145\ \mathrm{MeV})^4$. To take also the correction for the spurious collective motion of the centre-of-mass of the quark distribution relative to the cavity centre into account, we replace[6]

$$\Sigma\, \varepsilon_{\kappa n} \rightarrow [(\Sigma\varepsilon_{\kappa n})^2 - \Sigma\omega_{\kappa n}{}^2]^{\frac{1}{2}}, \qquad (11)$$

thus subtracting the contribution of the average total momentum from the total squared quark kinetic energy (this corresponds to the value $Z_o = 1.12$ for the constant in the conventional correction term[2] for the nucleon, which for giant bags becomes negligible anyway).

A giant quark bag in its ground state will adjust the numbers of up- and down-quarks (by a given nucleon number A) and the radius R such that the pressure of the quarks on the bag surface is balanced by the pressure exerted by the surrounding vacuum, or, equivalently[1], that the total energy E_{tot} is minimal:

$$M_{GQN}(A) = \min_{\{Z,R\}} E_{tot}(A,Z,R). \qquad (12)$$

This minimization will be performed in two steps. From

$$\partial/\partial R\ E_{tot}|_{R=R_o} = 0 \qquad (13)$$

we obtain immediately, if we neglect for a moment the finite mass of the down-quark and recognize the R-independence of both $\omega_{\kappa n}$ and ω_c,

$$R_o(A,Z) = \{[((\Sigma\omega_{\kappa n})^2 - \Sigma\omega_{\kappa n}{}^2)^{\frac{1}{2}} + \omega_c(Z)]/4\pi B\}^{\frac{1}{4}} \qquad (14)$$

for the equilibrium radius, and

$$E^o_{tot}(A,Z) = 16\pi BR_o{}^3(A,Z)/3 \qquad (15)$$

for the mass of the giant bag. From (14) and (15) it is clear that R_o and consequently also M_{GQN} will be minimal, if Z is adjusted to minimize the nominator in (14). To develop a convenient algorithm for the filling scheme including the effects of the Coulomb interaction we first list below two recurrence relations:

$$\Delta_u E_c(Z_u,Z_d) \equiv E_c(Z_u+1,Z_d) - E_c(Z_u,Z_d)$$

$$= 4\alpha/15R\ (2Z_u-Z_d) \equiv \Delta_u\omega_c(Z_u,Z_d)/R,$$

$$\Delta_d E_c(Z_u,Z_d) \equiv E_c(Z_u,Z_d+1) - E_c(Z_u,Z_d)$$

$$=-2\alpha/15R\ (2Z_u-Z_d) \equiv \Delta_d\omega_c(Z_u,Z_d)/R. \qquad (16)$$

As an example for their application we calculate the masses of the first four nuclei with A = 1,2,3,4. For this purpose only the two lowest quark levels have to be taken into consideration. The degeneracy of each level is $2j+1 = 2|\kappa|$ for angular momentum, 3 for colour and 2 for isospin. The first $(1s_{\frac{1}{2}})$ level has the eigenfrequency

$$\omega_{-1,1} = 2.043 \qquad (17)$$

and a total degeneracy of 12. The quark content of the A = 1 bag, i.e., the "nucleon", is (uud), if the stable lowest energy configuration is a proton, and (udd), if the stable configuration is a neutron[20]. In both cases however, one up and one down-quark are present. We therefore obtain first

$$RE_{ud} = \omega_{-1,1} + \Delta_u\omega_c(0,0) + \omega_{-1,1} + \Delta_d\omega_c(1,0) \qquad (18)$$

$$= 4.086 - 4\alpha/15,$$

and than for the proton configuration (udu)

$$RE_{udu} = RE_{ud} + \omega_{-1,1} + \Delta_u\omega_c(1,1) \qquad (19a)$$

$$= 6.129 - 4\alpha/15 + 4\alpha/15 = 6.129$$

and, respectively, for the neutron configuration (udd)

$$RE_{udd} = RE_{ud} + \omega_{-1,1} + \Delta_d\omega_c(1,1) \tag{19b}$$

$$= 6.129 - 4\alpha/15 - 2\alpha/15 = 6.129 - 2\alpha/5$$

Obviously, the neutron has lower mass than the proton (if the nucleon radius is 1 fm the corresponding energy difference is ~ 0.58 MeV) and is a stable particle in this approach with the down-quark mass kept at zero. Let us now proceed with constructing the next heavier nucleus. First we add again another ud-combination to the neutron to obtain

$$RE_{uddud} = RE_{udd} + 2\omega_{-1,1} + \Delta_u\omega_c(1,2) + \Delta_d\omega_c(2,2) \tag{20}$$

$$= 10.205 - 2\alpha/5 + 0 - 4\alpha/15 = 10.205 - 2\alpha/3,$$

and by inspection of the neutron-proton configuration

$$RE_{uddudu} = 12.248 - 2\alpha/3 + \Delta_u\omega_c(2,3) \tag{21a}$$

$$= 12.248 - 2\alpha/5$$

and of the di-neutron configuration

$$RE_{uddudd} = 12.248 - 2\alpha/3 + \Delta_d\omega_c(2,3) \tag{21b}$$

$$= 12.248 - 4\alpha/5$$

the di-neutron (21b) turns out to be the next stable configuration. If we add another three quarks we find the tri-neutron to again fullfil the condition of minimal energy. Since now the Pauli principle forbids to add another down-quark into the $1s_{\frac{1}{2}}$-shell, and the next eigenmode

$$\omega_{2,1} = 3.204 \tag{22}$$

is appreciably higher, the A=4-nucleus will contain one proton, because the addition of an upper-quark into the $1s_{\frac{1}{2}}$-shell is energetically favourable over adding a down-quark into the $2p_{3/2}$-shell. This action of the Pauli principle has apparently similar consequences as the "symmetry energy" in the conventional language of the liquid-drop theory.

For each case, i.e., for A=1,2,..., we thus chose the charge $Z = Z_{min}$, which corresponds to minimum Coulomb energy, and calculate the equilibrium radius $R_o(A)$ from eq. (14) to obtain the total bag mass

$$M_{tot}(A) = E^o_{tot}(A, Z=Z_{min}). \tag{23}$$

The equilibrium radius is plotted in Fig. 1 as function of the mass number A in the form

$$r_o(A) = R_o(A)/A^{1/3}. \tag{24}$$

The dotted curve is $r_o(A)$ calculated for bags containing only free quarks (no Coulomb interaction). As compared to the value 1.2 fm, known from conventional nuclear physics, the quark gas nuclei reach a rather significant compression. This compression is acted against by the Coulomb-force, i.e., $r_o(A)$ becomes greater when (8) has been included (dashed line). Finally, the full line has been obtained after the finite mass of the down-quark was incorporated, too. This is easily done by setting

$$\mu(A) = m_d R = m_d R_o(A) = m_d\ A^{1/3}\ 0.989\ \text{fm} \tag{25}$$

for the dimensionless mass, and performing the analysis as indicated above (the value of r_o has been chosen to be asymptotically, i.e. for large A, self-consistent, cf. Fig. 1). This obviously allows, and is justified at least for $A \geq 100$, to avoid complications in the minimalisation procedure due to the R-dependence of μ (now, moreover, the proton is also stable and the helium contains two protons).

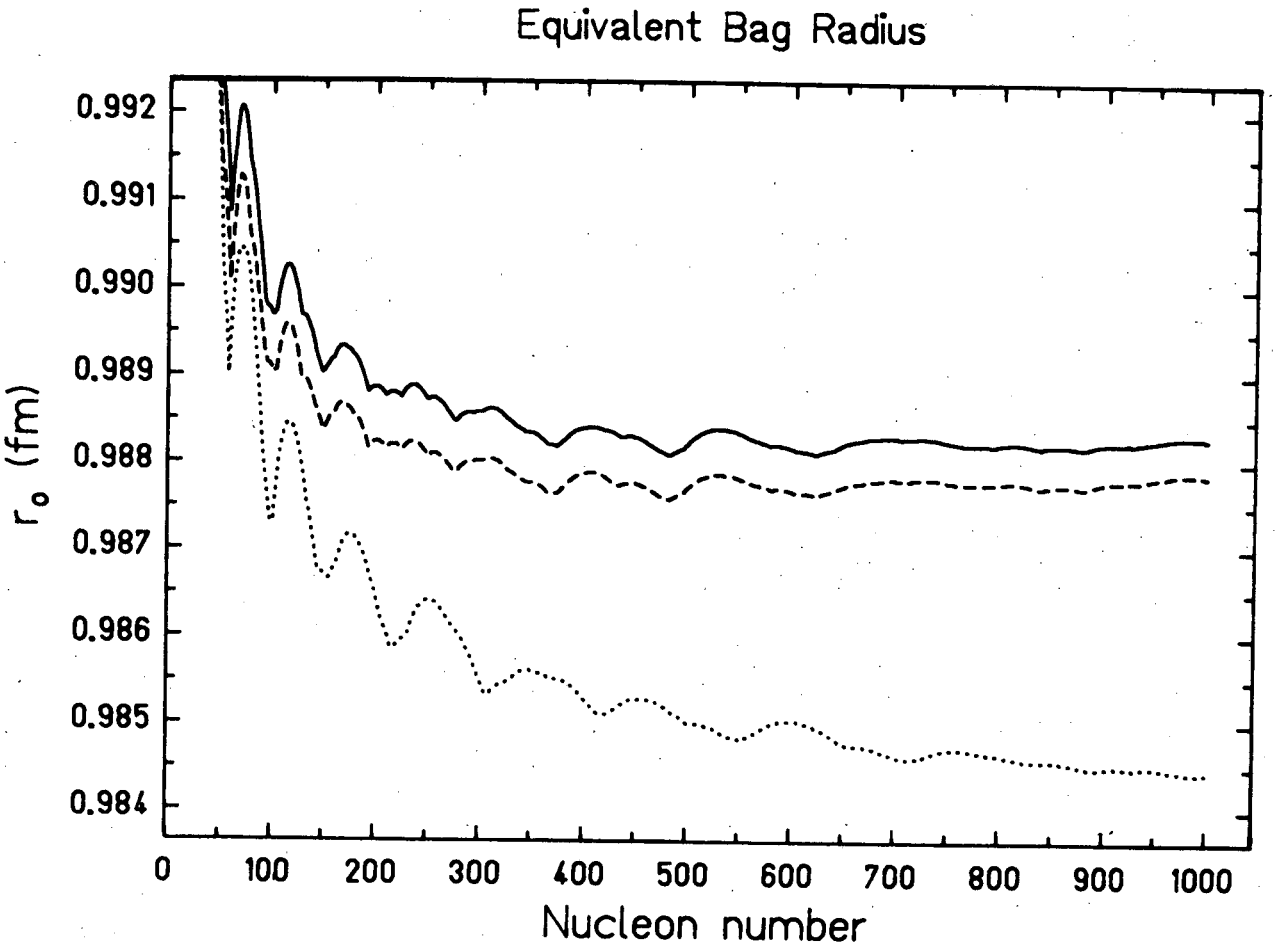

Figure 1. The equivalent radius of the GQN as function of A: Obviously these object are significantly compressed as compared with the ordinary nuclear matter (r_o of the order 1.1 - 1.2 fm). The compression is smaller, however, when the action of the Coulomb-force is allowed for (dashed line) and the down-quark mass is finite (full line). The structure of the curves reflects the shell-effects.

To illustrate the relative contribution of the shell-structure (symmetry energy), the Coulomb-energy and the finite quark-mass, the binding energies per nucleon

$$B.E.(A) \equiv -[\, M_{tot}(A)/A - m_{nucl}^{exp}] \tag{26}$$

with

$$m_{nucl}^{exp} = 938.3 \text{ MeV} \tag{27}$$

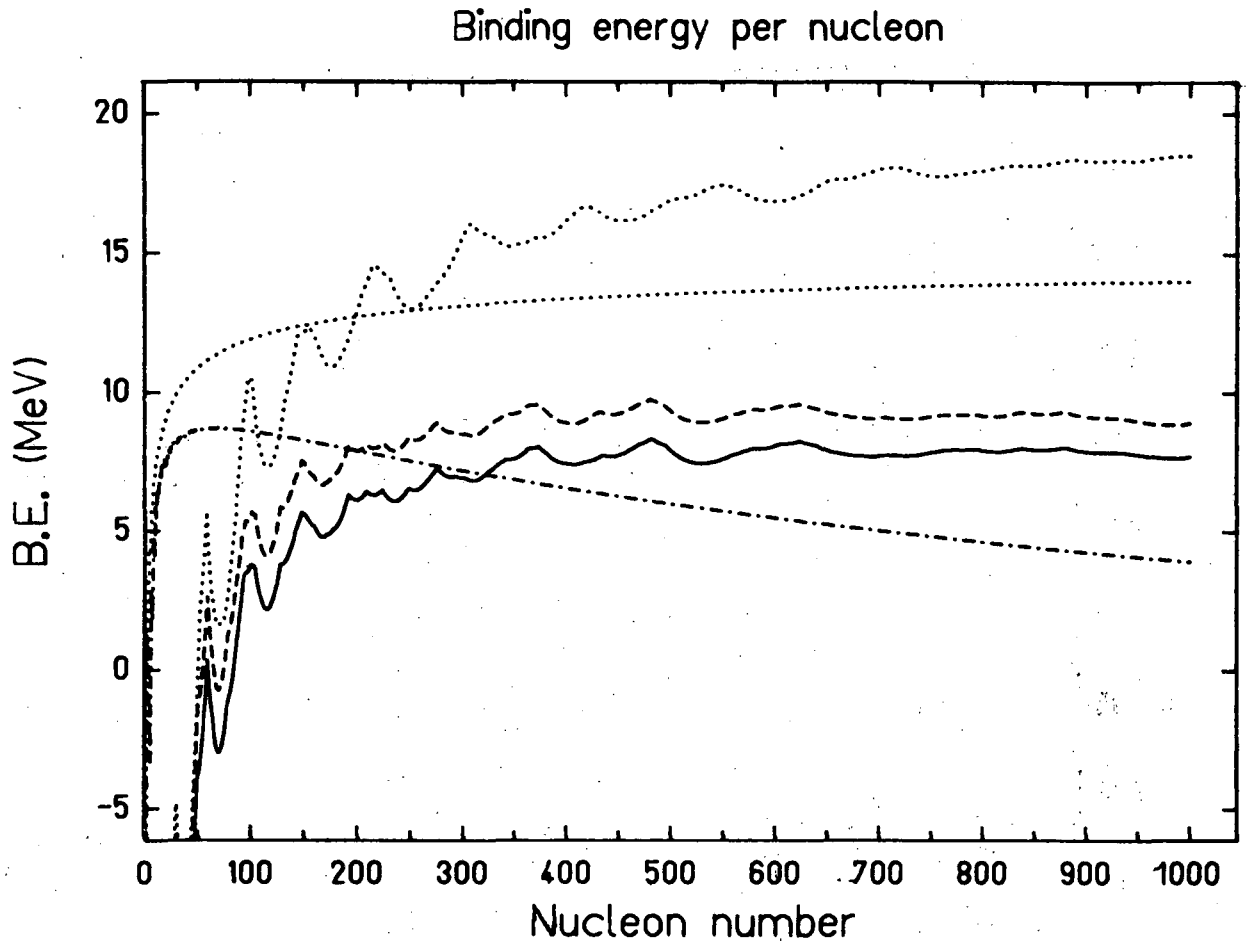

Figure 2. Binding energy per nucleon as function of A: If the Coulomb-energy is excluded the binding energy for giant quark nuclei (dotted curve) exceeds that of ordinary nuclear matter, discribed by the Bethe-Weizsäcker formula without the Coulomb-term (dotted smooth line), in the uranium region. The crossing point shifts, if the Coulomb energy is included in both formulas (Bethe-Weizsäcker without pairing, dash-dotted smooth line, and our bag-calculation, dashed curve) towards heavier nuclei. The shift is amplified, if the mass of the down-quark is nonzero (m_d = 5 MeV, full curve). In all cases the quarks do not interact via gluon-exchange in the bag-calculations.

are displayed in Fig. 2. In formula (26) we have subtracted the experimentally determined nucleon mass and ***not*** the "consistent" mass of the A=1-bag with free quarks

$$m_{nucl}^{cons} = M_{tot}(A=1) \approx 1217 \text{ MeV}. \tag{28}$$

Note, that otherwise the binding energy would be of the order of 300 MeV/A! The subtraction of the experimental proton mass can be justified in the following way. Although we assume that the residual colour interaction is negligible for hadrons with very large baryon numbers, we must not omit the one-gluon exchange in the case of hadrons with baryon number as small as 1! We know that in the MIT-bag-model colour hyperfine splitting resolves the N-Δ degeneracy, i.e. it brings the nucleon mass down to the experimental value and shifts the Δ-resonance ~ 300 MeV above it, using parameters consistent with ours[19]. The proton mass of 938.3 MeV is thus the value which in fact is the consistent one, provided our assumption, the residual interaction (i.e., in particular, the hyperfine splitting) be negligible for giant bags, is correct.

III. STRUCTURE OF GIANT QUARK NUCLEI.

In this Section the assumptions hidden in our model will be listed once more and their physical consequences will be carefully discussed. Basic to our considerations is the applicability of the MIT-bag model. This approach differs from other bag models in the way how the surface of the confining cavity is treated and how the chiral symmetry with its associated pion field is taken care of. The discontinuous surface used in the MIT-approach leads to severe conceptional difficulties[2] (for example the contradiction between the finite volume energy within and the infinite quark mass outside the bubble), but allows, on the other hand, for substantial simplification of the calculational effort with good "first order" results in explaining hadron spectroscopy[19]. A model in which the surface is treated properly in the form of a scalar field coupled to the fermion field is the soliton-bag model[4]. We expect, however, that with increasing number of quarks the inhomogenity in the scalar field, i.e. the "hole" in the vacuum-condensate digged by the quarks, will simply grow in its spatial extension, not changing the properties of the bag

walls, such that the contribution of the surface energy will decrease in relation to the volume energy, reaching thus the MIT-limit for very large baryon numbers.

In this sence we expect the MIT-approach to be applicable, but still have to be aware of the sensitivity of the results presented here with respect to the choice of the parameters, in particular of the bag pressure B which fixes the energy scale. Although the value B = (145 MeV)^4 was used in the original hadron fit[19], much has been argued against this choice[18] and fits to the hadronic mass spectrum were performed with modified values of B, but with other corrections included (for example the contribution of the pions[5], gluons[18], etc.). There exists, however, no unique prescription how these corrections should be accounted for[21], which graphs should contribute and which should be left out (self-energy!). Again, to be consistent with the neglect of the gluon exchange and with our aim to present a calculation based on as simple assumptions as possible, we stick to the conventional treatment of the MIT-bag and use the value of the bag pressure B = (145 MeV)^4. Note that since the total mass of the GQN is approximately proportional to $B^{\frac{1}{4}}$, a slight change of this constant will have a large effect onto the binding energy. For example, if the mass per nucleon of a GQN is 1 GeV, a 5% increase of $B^{\frac{1}{4}}$ will decrease the binding energy per nucleon by desastrous 50 MeV!

The second point to be mentioned is the homogeneous-sphere approximation for the evaluation of the Coulomb energy. In order to test its validity let us consider the exact quark distribution in the bag, which is given by

$$\rho_i(r) = \Sigma\, \psi_i^+ \psi_i \qquad i = u,d \tag{29}$$

where the sum runs over all occupied modes for up- (i=u) and down-quarks (i=d), and

$$\Psi_i^{nk\nu}(\vec{r}) = A_{kn} \begin{pmatrix} j_{\ell_k}\left(\frac{\omega_{kn} r}{R_0}\right) \chi_k^\nu(\Omega) \\ -\frac{i\,\omega_{kn}}{\varepsilon_{kn}+\mu}\, \mathrm{sgn}(k)\, j_{\bar{\ell}_k}\left(\frac{\omega_{kn} r}{R_0}\right) \chi_{-k}^n(\Omega) \end{pmatrix} \tag{30}$$

is the quark wave-function (note that $\mu=0$ and $\varepsilon=\omega$ for up-quarks) with the angular momentum

$$\ell_\kappa = \begin{cases} \kappa & \text{for} \quad \kappa>0 \\ -\kappa-1 & \quad\quad \kappa<0 \end{cases}, \quad \bar{\ell}_\kappa = \ell_{-\kappa} \tag{31}$$

and the magnetic quantum number

$$\nu = -j, -j+1, \ldots, j, \tag{32}$$

connected with the total spin

$$j = |\kappa| - \tfrac{1}{2}, \tag{33}$$

and the two component spherical spinors[22]

$$\chi_\kappa^\nu(\Omega) = \sum_m \sqrt{2j+1} \begin{pmatrix} \ell & \frac{1}{2} & j \\ \nu-m & m & -\nu \end{pmatrix} Y_\ell^{\nu-m}(\Omega)\chi^m \tag{34}$$

with

$$\chi^{\frac{1}{2}} = \begin{pmatrix} 1 \\ 0 \end{pmatrix}, \quad \chi^{-\frac{1}{2}} = \begin{pmatrix} 0 \\ 1 \end{pmatrix} \tag{35}$$

The normalisation factor is

$$A_{\kappa n} = \frac{|\omega_{\kappa n}|}{|j_{\ell_\kappa}(\omega_{\kappa n})| R_0^{3/2}} \left[2\varepsilon_{\kappa n}(\kappa+\varepsilon_{\kappa n}) + \mu\right]^{-\frac{1}{2}} \tag{36}$$

where $\varepsilon_{\kappa n}$ denotes the energy eigenvalue (1d) corresponding to the eigenspinors (30).

For closed shells with a specific value of κ we sum over all magnetic quantum numbers, obtaining

$$\sum_\nu (\Psi^{n\kappa\nu})^\dagger \Psi^{n\kappa\nu} = |A_{\kappa n}|^2 \left[j_{\ell_\kappa}^2 + j_{\bar{\ell}_\kappa}^2\right] \sum_\nu (\chi_\kappa^\nu)^\dagger \chi_\kappa^\nu \tag{37}$$

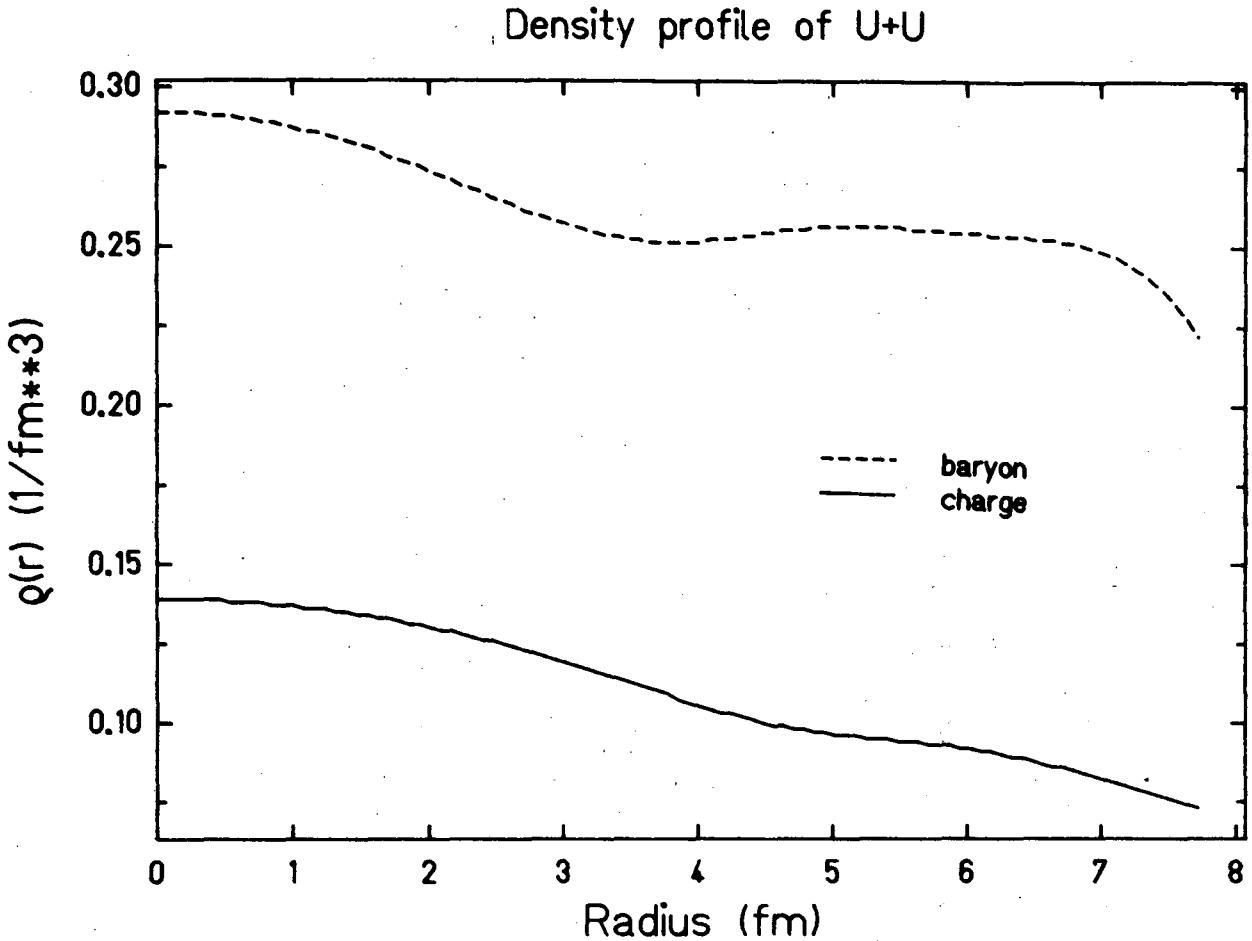

Figure 3. Density of double-uranium (A=476, Z=184): The dotted line is the baryon, and the full line the charge density distribution. This giant quark nucleus has a radius of 7.8 fm which should be contrasted with the radius of a conventional giant nucleus made of nucleons and bound to the standard law R = 1.2 fm $A^{1/3}$, i.e. R_{UU} = 9.4 fm.

The sum can be evaluated by inserting the definition (34) and profitting from the properties of the 3j-symbols[23]:

$$\sum_{\nu} (\chi_{\kappa}^{\nu})^{\dagger} \chi_{\kappa}^{\nu} = (2j+1) \sum_{\nu} \frac{1}{2\ell+1} |Y_{\ell}^{\nu}|^2 = \frac{2j+1}{4\pi} = \frac{2|\kappa|}{4\pi} \tag{38}$$

As expected, closed shells are rotationally invariant. For open shells this will not be the case; such quark nuclei can possibly deform, as do ordinary open shell nuclei. We neglect, however, the deformation of the bag also in this case, and replace the total degeneracy $2|\kappa|$ in (38) by the

number which corresponds to the actually occupied states in the open shell. The result is

$$\rho_i(r) = \sum_{\kappa n} \frac{1}{4\pi} D(\kappa) |A_{\kappa n}|^2 \left[j^2_{\ell_\kappa}\left(\frac{\omega_{\kappa n} r}{R_0}\right) + j^2_{\bar{\ell}_\kappa}\left(\frac{\omega_{\kappa n} r}{R_0}\right)\right] \tag{39}$$

with n and κ running now over occupied shells and $D(\kappa)$ being the corresponding quark number (i.e., $D(\kappa)=6|\kappa|$ for each isospin direction, if the shell is closed). Having calculated the expressions (39) for up- and down-quarks, the baryon density can be written

$$\rho_B = \frac{1}{3}\left(\rho_u + \rho_d\right) \tag{40}$$

and the charge density is

$$\rho_c = \frac{2}{3}\rho_u - \frac{1}{3}\rho_d \tag{41}$$

In Fig. 3 the baryon and charge densities for double-uranium are plotted as the physical quantities of interest (recall that from (40) and (41) $\rho_u = \rho_B + \rho_C$ and $\rho_d = 2\rho_B - \rho_C$). On the average we indeed find these nuclei more or less homogeneous, such that our approximation is expected to be correct at least on the qualitative level. However, the repulsion of the charge distribution towards the surface (polarisation of GQN) is not contained in this treatment, since the energy levels are not allowed to rearrange due to the Coulomb force, and a depletion of levels with low angular momentum (wave function concentrated in the centre of the bag) in favour of high angular momentum states (wave function concentrated towards the bag walls) does not occur in this zero'th order estimate. Obviously, a self-consistent (Hartree-Fock) calculation would be required to take the polarisation effects due to the Coulomb and eventually also colour forces into account. Clearly, much work is still to be done before reliable conclusions can be drawn.

In spite of this we shall try to glance at the physical implications of our model, and return to Fig. 2 again. The Coulomb-energy and the finite

quark-mass are not included in the upper (dotted) line displaying the binding energy per nucleon (i.e., per three quarks) which crosses the binding energies per nucleon obtained by extrapolating the empirically determined Bethe-Weizsäcker formula supposed to represent ordinary nuclear matter. The latter is given by[24]

$$B.E.(A,Z) = a_V - a_S A^{-1/3} - a_C Z^2 A^{-4/3} - a_A [(Z-\tfrac{1}{2}A)/A]^2 + a_P A^{-3/2}\,\delta, \qquad (42)$$

with the parameters

a_V = 15.85 MeV,

a_S = 18.34 MeV,

a_C = 0.71 MeV,

a_A = 92.86 MeV,

a_P = 11.46 MeV,

$$\delta = \begin{cases} +1 \text{ for even-even nuclei} \\ -1 \text{ for odd-odd nuclei} \\ 0 \text{ otherwise} \end{cases}$$

fitted to experimental data. The smooth dotted curve represents the Bethe-Weizsäcker formula with the Coulomb-energy omitted. It crosses the analogous curve for quark nuclei in the region of uranium.

In this case the Z/A-ratio is $\frac{1}{2}$ on average due to the action of the Pauli-principle. If we include the Coulomb interaction in the schematic way discussed previously, the binding energy (dashed curve) and the Z/A-ratio decrease. The Bethe-Weizsäcker formula with the Coulomb term included now, yields binding energies indicated by the smooth dash-dotted line. The crossing between the two curves shifts, however, only gradually. A much larger shift of the crossing-point into the region of giant elements

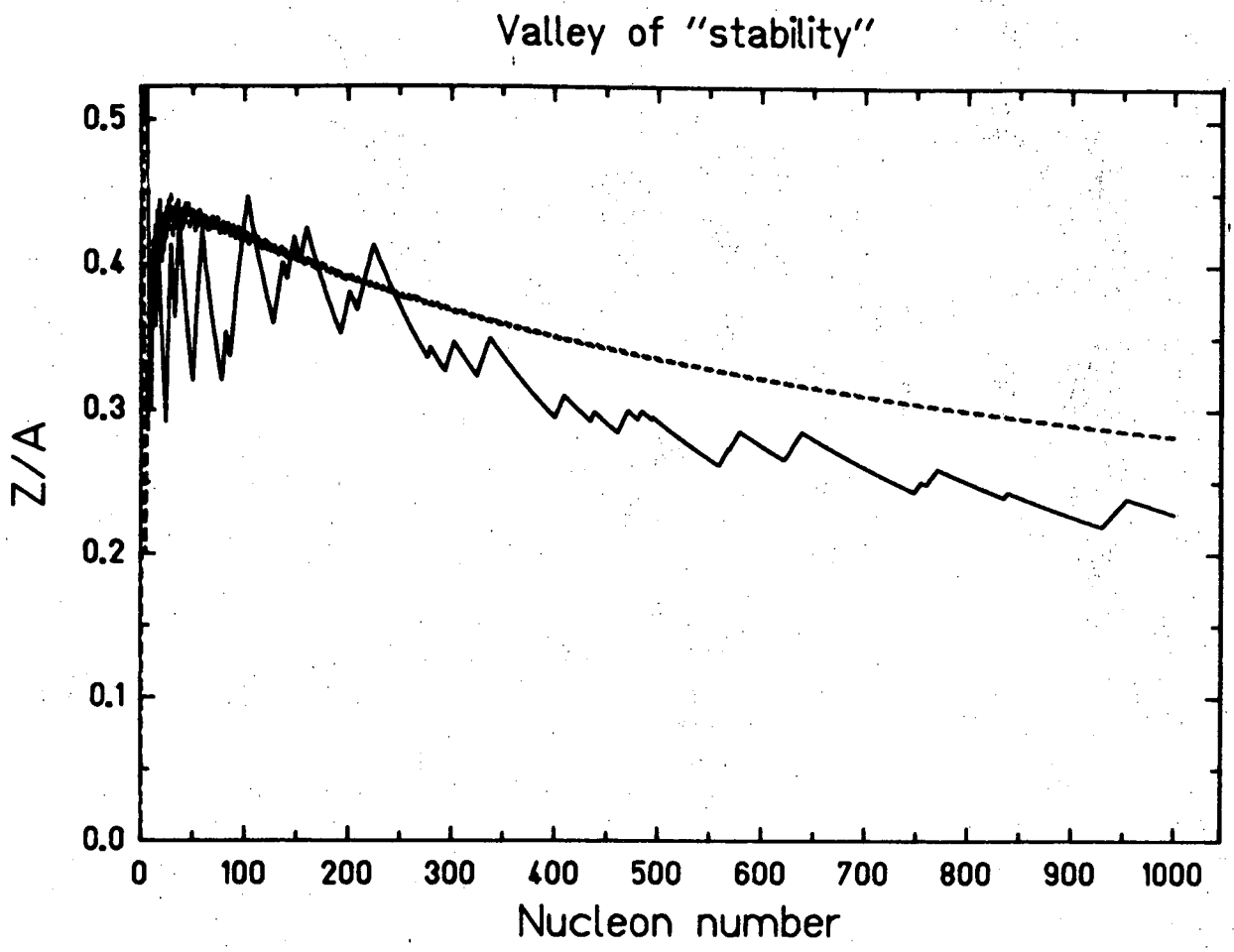

Figure 4. The Z/A-ratio: it is obtained by minimalisation of the binding energy with respect to the total charge Z with A held constant. The results of our bag-calculation (full line) differ significantly from those obtained in an independent minimalisation of the Bethe-Weizsäcker binding energy (dashed line) for high baryon numbers.

arises, if the finite mass of the down-quark (lower full line) is included in the determination of the masses of GQN. Note that on the average the Z/A-ratios obtained from the minimalisation of the Bethe-Weizsäcker formula (which is based on the liquid-drop picture for the atomic nucleus and does therefore not contain shell corrections) differ from those obtained in our bag-calculation (see Fig. 4) significantly in the region of giant nuclei. Giant quark nuclei tend, at least in our approximation, to large "neutron" excess. It is only the stable valleys ($Z = Z_{st}(A)$) for nucleonic and quark nuclei which is exhibited in Fig. 4. The "neutralisation" of a GQN with $Z > Z_{st}$ charges can proceed either via proton-, α-particle

emission or via e^+-decay. These decay modes of GQN deserve further investigation which we shall not carry out here.

The prediction of our model is obviously now the following: if clusters of nuclear matter with very large baryon numbers of the order of ~400 and more are created, the nucleons seize to exist and the interiour of these objects turns into a (more or less free) quark gas. This transition is expected to be very fast, i.e. of the typical nuclear time scale $\tau \leq 10^{-21}$ sec. Again, its time development is an important task for the future. The shell structure in GQN is different from the shell structure in ordinary nuclei, although, as noted already some time ago (see K. Bleuler et al., Ref. 12), the use of the Dirac equation (1) automaticaly generates the right ordering of angular momenta which in the non-relativistic oscillator model has to be enforced by including the spin-orbit coupling[25]. Magic numbers in the MIT-bag model are, however, 2, 6, 8, 14, (18,20), 28, 34,..., as compared with the conventional 2, (6), 8, (14,16), 20, 28, (38,40), 50, ... (in parentheses the minor shell closures are indicated). The differences are caused by shifts in the level ordering still present (e.g., the 2s-level (corresponding to the Dirac quantum number $\kappa=-1$) is between the 1d-levels ($\kappa=-3$ and 2, respectively) in the nuclear shell model, whereas it lies above them in the MIT-bag and is immediately followed by the first f-level ($\kappa=-4$), cf. Table I).

Now, of course, the question arises, why no experimental hints for such anomalous objects have so far been reported. One has to realize, however, that the search for anomalous states of nuclear matter has been put forward in high-energy heavy-ion collisions which are able to reach high-temperatures and high-density phases of nuclear matter due to the existence of nuclear shock waves[26], but are limited to small projectiles. Effects of large baryon numbers can be seen only with machines which are able to accelerate projectiles up to the heaviest ones. Since such devices like the UNILAC at GSI are, on the other hand, yet limited with respect to the ion final velocity, nobody has looked for such effects, particularly because of the common expectation that quark structure is restricted to show up in the high-energy domain.

System	A	Z
Pb + Pb	416	164
Ta + U	419	165
Au + U	435	171
Pb + U	446	174
Pb + Cm	455	178
Th + Th	464	180
Th + U	470	182
U + U	476	184
Cm + Th	479	186
Cm + U	485	188

Table II: Baryon and charge numbers of several realistic "giant" collision systems presently under investigation at GSI.

As already mentioned above, an anomalous positron production has been measured in recently reported experiments with heavy ions accelerated up to velocities in the vicinity of the Coulomb-barrier. These experiments were carried out over nearly a decade to study the expected instability of the QED-vacuum in very strong ("supercritical") electric fields generated by giant nuclear systems[14]. For long-living such systems ($\tau \geq 3\times10^{-20}$ sec) a sharp resonance should appear in the positron spectrum, when the overcritical charge of the united system is partially screened by electrons created out of the vacuum, and the total charge is balanced by positron emission. Such positron lines have indeed been found in various experiments[27], but its dependence on the positron kinetic energy seems not to show the scaling with $Z = Z_{proj} + Z_{target}$, expected on the grounds of standard extrapolation of the properties of ordinary nuclear matter to the domain of giant nuclei. The position of the "spontaneous" peak in the positron spectrum is rather energy independent at $T_{pos} \sim 320$ keV.

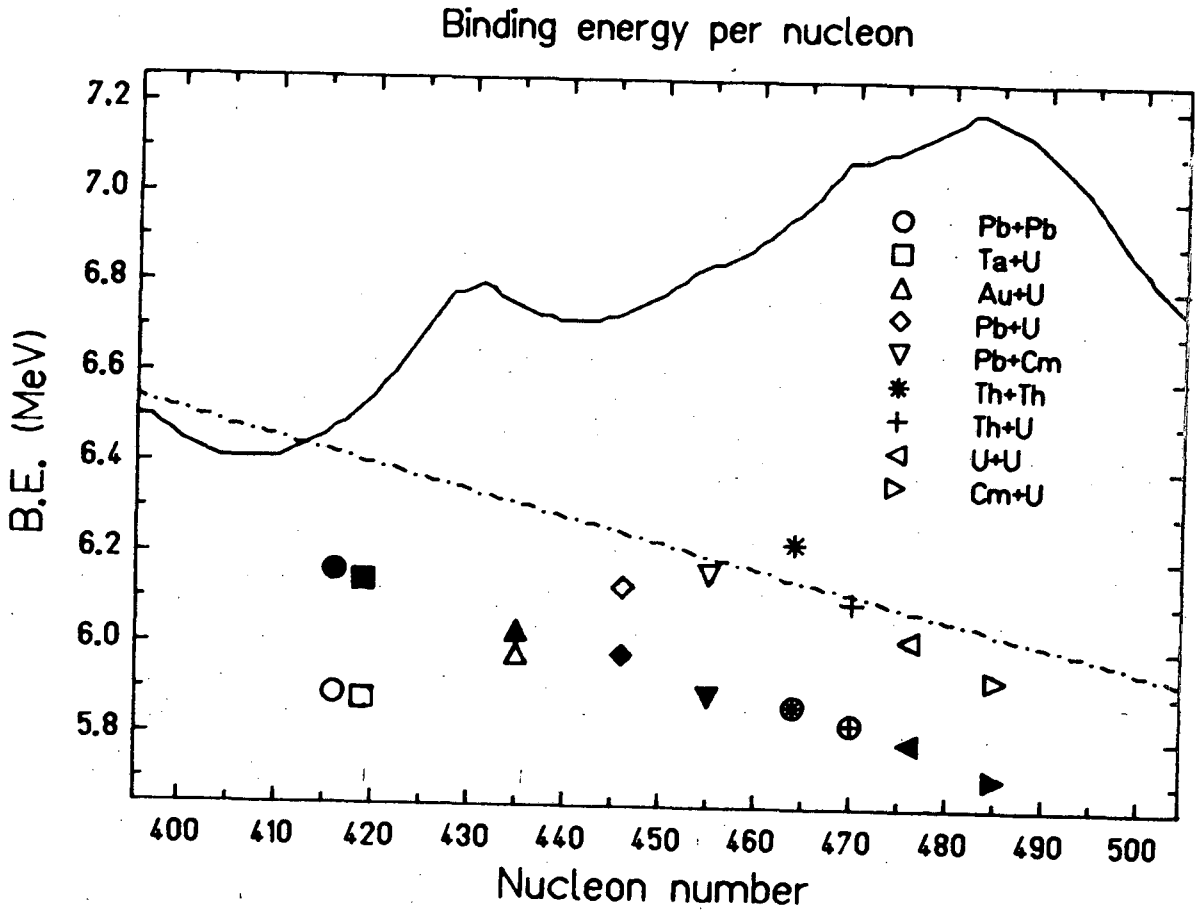

Figure 5. Binding energy of realistic giant quark systems: The open symbols are the GQN, the solid ones the corresponding normal nuclei (Bethe-Weizsäcker formula). The solid line is the one already displayed in Fig. 2, the dash-dotted line is Bethe-Weizsäcker matter.

What happens now, if we suppose that giant quark nuclei are created in these experiments? Can the existence of a GQN account for these observations? Of course, the nuclear physics, which so sensitively determines the above mentioned scaling behaviour of the positron resonance energy, is dramatically changed. However, in trying to answer this question we have to be careful, since the reaction systems with high baryon numbers used in actual experiments do not have a total charge wich coincides with the minimum charge Z as calculated in Section II. Their charge is rather fixed by the incoming more or less stable nuclei (e.g., U+U, U+Cm, U+Th, Th+Th, etc.), and since the expected life-time of the unified system is too short for weak processes, it remains unchanged (up to a possible emission of charged light clusters[28]) during the reaction.

System	B.E. (MeV)	EBW (MeV)	ΔB.E. (MeV)	R (fm)	T_{pos} (KeV)
Pb + Pb	5.89	6.16	-0.27	7.38	—
Ta + U	5.88	6.14	-0.26	7.39	—
Au + U	5.98	6.03	-0.05	7.49	22.48
Pb + U	6.13	5.98	+0.15	7.55	130.65
Pb + Cm	6.16	5.89	+0.27	7.60	291.81
Th + Th	6.23	5.87	+0.36	7.65	376.66
Th + U	6.10	5.83	+0.27	7.69	467.73
U + U	6.02	5.79	+0.23	7.72	562.10
Cm + Th	5.96	5.75	+0.21	7.73	663.31
Cm + U	5.93	5.71	+0.22	7.77	767.78

Table III: The binding energies (B.E.) of GQN, equally sized normal nucleus (EBW), and the gain in binding energy (ΔB.E.) after transition into the quark gas state. The radius of the GQN as well as the positron kinetic energies are also given.

In Table II the systems which have been or will be used in actual GSI-experiments are listed. The binding energies of these systems are not contained in Fig. 2, where only the nuclei along the valley of stability (Z minimised for given A) are displayed. And, since in the Bethe-Weizsäcker formula (without shell corrections) the charge has also been minimised to obtain the binding energy of ordinary nuclear matter in Fig. 2, we also have to compare our results for the combined systems with the binding energies of normal nuclei with the fixed charges, off the corresponding stable valley. This is listed in Table III and shown also in Fig. 5.

Note first that the difference in binding energy per nucleon for quark nuclei and ordinary nuclei is small in the region of baryon numbers between 400 and 500; in this calculation, where shell corrections for the liquid drop theory are omitted, ordinary nuclei seem to be favoured for systems like Pb+Pb, Ta+U, but quark nuclei particularly for Th+Th and its neighbours. These borders are, however, so sensitive to the parameters and details of the model that they should not be taken literally, but rather as a qualitative hint that a transitions between the two phases of nuclear matter can easily happen as a function of A (N and Z) and, of course, also as a function of excitation energy.

Obviously, the Th-Th-system appears to be more bound than the other systems under consideration. This is a consequence of the shell structure: for Th+Th the number of up-quarks is 644, the number of down-quarks 748 (cf. Table II and eq. (5)), both just two units above the respective "magic" numbers (642) ($4s_{\frac{1}{2}}$-shell filled) and 744 (first $\kappa=8$ shell filled). The energy gain of ~0.36 MeV/A relatively to the normal nuclear matter substantially lowers the repulsive barrier between the thorium ions, namely by $2\times A_{Th}\times 0.36$ MeV $\approx$ 167 MeV, and could be a new explanation of the long life-time of the unified system, sufficient to destroy any vacancy in the electronic 1s-level by spontaneous positron production. The arising decrease of the nucleus-nucleus potential at the spherical point is even, although a little too large, of the correct order of magnitude required to reach a very stable nuclear structure of these giant objects[14]. Moreover, since Th+Th seems to be the combination whith the tightest binding and hence with the possibly longest life-time, it is not off-hand that all spontaneous positrons are emitted from this system, even though the reaction systems are heavier. The missing mass is perhaps emitted in the form of small yet undetected fragments.

IV. OUTLOOK

If the positrons are always emitted from the same giant unified system, there is no miracle about their constant energy. On the other hand, the quark-gas Th+Th-system has a radius of ~7.66 fm, i.e., it is significantly smaller (similarly as also the other GQN are very much more compressed then their "brothers" made of nucleons) then the conventional unified Th+Th spherical giant system (radius of ~9.2 fm). Since the positron kinetic energy depends very sensitively upon the radius of the charge distribution, then, if taken seriously, the anomalous positron peak should[29] be centered around ~0.5 MeV in this case. This is, however, far above the experimentally found value of 320 keV. This fact as well as the too large binding energy mentioned in the last section indicate that we have perhaps overestimated the compression of GQN. A more careful theoretical analysis of the Coulomb- and colour corrections might improve on this point, though.

If the assertion that the positrons are always emitted from the same giant system seems inacceptable, then the position of the positron peak will, of course, depend on the collision system under consideration. Discarding for a moment all doubts we still have about the correctness of our results in view of their sensitivity to the choice of the model parameters and of the crude approximations made, the positron energies can easily be calculated[29]. The result is listed in Table III.

It also should be mentioned that since GQN are supposed to be highly polarizable, in particular by the action of the long range electromagnetic force, the arising strong electric and magnetic fields may, depending on the geometry of the system, induce significant enhancements of various types of fermionic (scalar, pseudoscalar, etc.) densities which in turn could serve as anomalous sources for ***non-perturbative creation of new particles***. Such particles, if decaying into the electron-positron channel, could eventually also produce sharp monoenergetic positron resonances[30].

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