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THE SWELLING HADRONS[†]

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ABSTRACT

The notion of a "swelled world" for strong interactions is introduced, followed by a discussion on some phenomenological consequences of the "dropping" meson and baryon masses in dense and/or hot nuclear matter.

1 Introduction

In QCD, the mass of a hadron, say, a proton is of the form

$$m_P = M_0 + \langle p|(m_u\bar{u}u + m_d\bar{d}d + m_s\bar{s}s)|p\rangle \tag{1}$$

where m_q is the current quark mass of flavor $q = u, d, s$ and M_0 is a constant independent of the flavor content of the baryons. It is generally believed that M_0 arises dynamically from spontaneous breaking of chiral $SU(N_f) \times SU(N_f)$ symmetry. In the chiral limit $m_q \to 0$, therefore, the hadron mass is given entirely by the dynamically generated one.

The mechanism for the mass generation M_0 is not well understood at the moment and one would like to get some idea how and when the mass appears and disappears. On a fundamental level where the unification of "everything" is involved, there is a deep question of where the masses in general come from: One would like to understand at the same time where the electron and, say, a top quark with a vastly differently mass scale get their masses. The particle physicists address this question by asking how those fundamental masses are generated as the primordial super-hot and super-dense universe expands and cools. At the level of nuclear physics that we are concerned with, one can ask a similar question:

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Knowing that the masses of nucleons and mesons that we measure in laboratories are generated "spontaneously" what does high temperature or high density do to the masses? These two questions are clearly not unrelated.

I will discuss in this talk possible phenomenological consequences of the new idea proposed recently by Gerry Brown and myself in [1] that as one increases temperature and/or density as in relativistic heavy-ion collisions or in stellar collapse, the hadron masses scale as

$$m_B^*/m_B \approx m_M^*/m_M \approx f_\pi^*/f_\pi \equiv \Phi \tag{2}$$

where the subscripts B and M stand respectively for baryons and mesons (other than Goldstone bosons), the superscript * denotes quantities at nonzero temperature T and/or density ρ [‡] and f_π is the pion-decay constant standing for the only scale parameter relevant for the problem. Strictly from the theoretical point of view, our scaling scenario is not entirely unique. Different inputs in the theoretical model could make different predictions [2] and it will ultimately be up to experiments to prove or disprove the scaling of the type (2). In this talk, I will limit myself to the scaling (2) which might be called "Brown-Rho" scaling to avoid confusion.

2 Hadron Scale

QCD in chiral limit, while classically scale- (or conformal-) invariant, gets a scale by quantum effects. The vacuum breaks chiral symmetry, developing a non-vanishing quark condensate $\langle\bar{q}q\rangle$ which in turn endows masses to the hadrons and generating Nambu-Goldstone bosons (Goldstone bosons in short). This gives a scale, characterized by nonvanishing pion decay constant f_π. Quantum field theory, to be well-defined, requires regularization that involves a dimensionful cut-off Λ, thereby bringing in a scale. QCD, when quantized, has therefore a nonvanishing trace of energy-momentum tensor θ_μ^μ – known as trace anomaly – of the form

$$\theta_\mu^\mu = -\frac{\beta}{g}\,\mathrm{Tr}\; G_{\mu\nu}G^{\mu\nu} \tag{3}$$

where β is the renormalization group β function, g the color gauge coupling and G the gluon field tensor. That the trace of the energy-momentum tensor is non-vanishing is equivalent to saying that the divergence of the dilatation current is nonvanishing,

$$\partial^\mu D_\mu \neq 0, \qquad D_\mu \equiv x^\nu \theta_{\mu\nu}, \tag{4}$$

that is to say, conformal invariance is broken quantum mechanically. The scale so generated implies a nonvanishing gluon condensate $\langle\, \mathrm{Tr}\; G_{\mu\nu}G^{\mu\nu}\rangle$.

‡From now on unless otherwise specified, I will concentrate on density effect. Similar discussions can be made with temperature and I will make references to some of them.

Since the hadron masses are generated dynamically, they must be functions of the condensates, say,

$$M = f\left[\langle\bar{q}q\rangle, \langle \text{Tr } G_{\mu\nu}G^{\mu\nu}\rangle\right]. \tag{5}$$

We cannot calculate the precise form with which the condensates appear in the masses but we can calculate the ratio Φ. This we will do in what follows using effective chiral Lagrangians.

3 The In-Medium Effective Lagrangian

The key argument of [1] can be summarized as follows. The idea of effective Lagrangians is to get as close as possible to physical observables at the *tree level* of the Lagrangian with loop corrections suppressed by the cut-off scale mass. Here we are concerned with chiral invariance and scale invariance, so the construction of our effective Lagrangian can be streamlined along these properties of QCD. In general, implementing chiral and scale invariance in an effective Lagrangian is not unique. To make our argument predictive, we make one basic assumption which appears to be reasonable. We assume that at any density $\rho^* \leq \rho_{critical}$ where $\rho_{critical}$ is some critical density at which a phase transition – be that meson condensation or chiral transition – is to take place, *the effective Lagrangian continues to respect the same approximate chiral and scale invariances.* At the moment, we have no strong theoretical argument why this should be so but we will see later that nature seems to support this hypothesis. Whether or not this can be pushed all the way up to the QCD chiral phase transition is not clear. Nonetheless a rather compelling argument for such an extrapolation has been put forward in connection with QCD phase transitions [3].

For the moment, we will continue ignoring current quark masses. The explicit symmetry breaking will be brought in later. In the chiral limit, the vacuum expectation value of θ_μ^μ involves only the gluon fields as one can see in (3). Since the gluon field has canonical scale dimension -1, θ_μ^μ has scale dimension -4, so we may define [4] a scalar glueball (or dilaton) field χ of scale dimension -1,

$$\theta_\mu^\mu \equiv \chi^4 \tag{6}$$

and express the gluon condensate as $\langle 0^*|\chi^4|0^*\rangle$. This scalar glueball field plays an important role in implementing trace anomaly to effective chiral Lagrangians.

In what follows, I will discuss in terms of only the scalar field χ and the chiral field $U = e^{i\pi/f_\pi}$ but one can easily generalize the consideration to vector-meson and baryon fields using arguments based, respectively, on hidden gauge symmetry and skyrmion description. For simplicity, I will restrict to the $SU(2)$ flavor. The generalization to the $SU(3)$ flavor is straightforward.

Chiral invariance requires that there be a quadratic current algebra term

$$\sim \text{Tr}(\partial_\mu U\partial^\mu U^\dagger),$$

a quartic (Skyrme) term,

$$\sim \mathrm{Tr}[U^\dagger \partial_\mu U, U^\dagger \partial_\nu U]^2$$

and so on. These terms, not involving glueball fields, should not be expected to contribute to the trace of the energy momentum tensor. Therefore we would like them to be scale invariant. Since the chiral field U has scale dimension 0 because of its non-linear realization of chiral symmetry [5], the quartic term clearly has scale dimension −4, so it is scale-invariant. But the quadratic term has scale dimension −2 and hence does not have the right dimension. We can remedy this by multiplying it by an object of scale dimension −2, *i.e.*, by χ^2. We can continue doing this to *all* the chiral derivative terms by simply using various powers of the scalar glueball field. Now having assured that the chiral fields do not contribute to θ^μ_μ, we need to have a term made up of the glueball fields that will supply the contribution $\sim \chi^4$. We may do this by adding a potential term [4]

$$V(\chi) \sim \chi^4 \ln \chi. \tag{7}$$

Now what about the explicit chiral symmetry breaking which introduces an additional scale parameter? This is an intricate story which introduces a certain element of uncertainty in constructing effective Lagrangians. The reason is that the quark mass gives an additional term to θ^μ_μ with an anomalous dimension contribution γ (because of the chiral symmetry breaking by the mass term)

$$\theta^\mu_\mu = \sum_q (1+\gamma_q) m_q \bar{q} q - \frac{\beta}{g}\, \mathrm{Tr} G_{\mu\nu} G^{\mu\nu}. \tag{8}$$

Apart from the anomalous dimension contribution, the mass term brings in a dimension −3 term. How does one represent such a term in terms of the chiral fields? Here I will take the most obvious possibility which is to multiply by a scale dimension-3 field χ^3 to the mass term $\mathrm{Tr}(MU + h.c.)$ (of scale dimension 0) where M is the quark mass matrix. Just as the anomalous dimension modifies the scale dimension of the quark mass term in the QCD Lagrangian, we would expect that quantum effects *with a given effective Lagrangian* would bring in terms of other scale dimensions. While the anomalous dimension in QCD proper may be ignorable, it is possible that loop effects in effective theories may not be small *for Goldstone bosons*. Indeed a reason is offered later as to why loop effects are expected to be important for pion mass –and perhaps for kaon mass as well.

Collecting all the relevant terms together, our effective Lagrangian is of the form

$$\begin{aligned}\mathcal{L}^{eff} = \frac{f_\pi^2}{4}\left(\frac{\chi}{\chi_0}\right)^2 \mathrm{Tr}(\partial_\mu U \partial^\mu U^\dagger) + \frac{1}{32e^2}\, \mathrm{Tr}[U^\dagger \partial_\mu U, U^\dagger \partial_\nu U]^2 + \ldots \\ + c\left(\frac{\chi}{\chi_0}\right)^3 \mathrm{Tr}(MU + h.c.) + \ldots \\ + \frac{1}{2}\partial_\mu \chi \partial^\mu \chi + V(\chi).\end{aligned} \tag{9}$$

where the ellipsis stands for higher derivative and mass matrix terms that can contribute. I have added the scale-invariant kinetic energy term for the scalar field. The quantity χ_0 is a constant of mass dimension which will be identified with the vacuum expectation value of the χ field in medium-free space, $\chi_0 \equiv \langle 0|\chi|0\rangle$.

The potential $V(\chi)$ is manufactured such that it gives the trace anomaly correctly [4]. One can always add scale-invariant and chiral-invariant quantities to it. We will assume that for a given density ρ^*, the potential is minimized at a vacuum expectation of the χ field,

$$\chi_* \equiv \langle 0^*|\chi|0^*\rangle. \tag{10}$$

This suggests to expand the Lagrangian around the vacuum value by shifting

$$\chi = \chi_* + \chi' \tag{11}$$

where χ' is the fluctuating scalar field. Define

$$f_\pi^* \equiv f_\pi \frac{\chi_*}{\chi_0}. \tag{12}$$

The Lagrangian (9) now reads

$$\begin{aligned}\mathcal{L}^{eff} = \frac{f_\pi^{*2}}{4}\,\mathrm{Tr}(\partial_\mu U\partial^\mu U^\dagger) + \frac{1}{32e^2}\,\mathrm{Tr}[U^\dagger\partial_\mu U, U^\dagger\partial_\nu U]^2 + \ldots \\ +c\left(\frac{f_\pi^*}{f_\pi}\right)^3 \mathrm{Tr}(MU + h.c.) + \ldots \\ +\frac{1}{2}\partial_\mu\chi\partial^\mu\chi - \frac{1}{2}m_\chi^{*2}\chi^2 + \ldots,\end{aligned} \tag{13}$$

with

$$m_\chi^{*2} = m_\chi{}^2\left(\frac{\chi_*}{\chi_0}\right)^2 \approx m_\chi^2\left(\frac{f_\pi^*}{f_\pi}\right)^2. \tag{14}$$

The last ellipsis in this equation stands for all possible χ field couplings to other fields that enter into the Lagrangian.

We can now readily extract the consequences of this procedure for the *tree level* masses and coupling constants. Baryons emerge as skyrmions of this Lagrangian and hence their masses scale as

$$m_B^* \sim f_\pi^*/e \sim f_\pi^*\sqrt{g_A^*} \tag{15}$$

from which we obtain

$$\frac{m_B^*}{m_B} \approx \frac{f_\pi^*}{f_\pi} \equiv \Phi(\rho), \tag{16}$$

while

$$\frac{g_A^*}{g_A} \approx 1. \tag{17}$$

The in-medium Goldberger-Treiman relation implies that the πNN coupling scales as $\frac{g^*_{\pi NN}}{g_{\pi NN}} \approx 1$. Vector mesons emerge as hidden gauge bosons from the chiral Lagrangian [6] from which follows the mass formula

$$m_V^{*\,2} = 2f_\pi^{*2}e^2 \tag{18}$$

which leads to

$$\frac{m_V^*}{m_V} \approx \Phi(\rho). \tag{19}$$

To the leading order in N_c, this scaling applies to the ρ, ω, A_1 and other vectors. Finally we expect that the quarkish scalar meson σ that figures in nuclear physics as an interpolating field for a two-pion correlation, coupling with the scalar glueball [7], will scale as

$$\frac{m_\sigma^*}{m_\sigma} \approx \frac{m_\chi^*}{m_\chi} \approx \Phi. \tag{20}$$

Thus we have a "universal scaling" characterized by one parameter, Φ. But this applies to those hadrons whose masses are mainly generated by spontaneous symmetry breaking. In the chiral limit, Goldstone-boson masses are zero and remain zero at all density, as they are protected by chiral invariance. Therefore the behavior of the pion or kaon masses must reflect directly on the way the chiral symmetry is *explicitly* broken at the electroweak scale. This implies that how their masses change in density or temperature must be quite intricate.

Be that as it may, the effective Lagrangian (13) predicts at tree level that the pion mass scales

$$\frac{m_\pi^*}{m_\pi} \approx \sqrt{\Phi}. \tag{21}$$

But this is not the entire story; it is not even a dominant part of it. The reason is simply that there are other terms which are equally important. For instance, Pauli-blocking effects, appearing at one-loop level are enough to compensate the scaling of (21). The intricacy associated with the pion mass is described elsewhere [8]. The upshot of the story is that the pion mass does not seem to change in any significant way up to the density we are concerned with. We will simply take it that the pion mass does not change at all. (See [8] what happens to the pion mass in hot and dense system.)

4 Chiral Perturbation Theory in Nuclei

Given the effective Lagrangian (13), how do we go about seeing the effect of dropping masses in nuclei? My proposal is that we do a chiral perturbation calculation with the Lagrangian following a recent proposal by Weinberg [9]. The only subtle point here is that the masses and the pion decay constant are density-dependent.

Weinberg has given a dimensional argument to show that for processes involving small momentum Q, an amplitude $\mathcal{A}$ with incoming or outgoing n_π number of pions and n_e number of nucleons can be characterized by

$$\mathcal{A} \sim Q^\nu f(Q/\mu), \tag{22}$$

$$\nu = 2 - \frac{1}{2}n_e + 2L + \sum_i V_i[d_i + \frac{1}{2}n_i - 2] \tag{23}$$

with μ a scale parameter of order $\Lambda_\chi \sim 1$ GeV, n_e the number of external nucleon lines, L the number of loops, V_i the number of vertices of the type i, d_i the number of derivatives acting on ith vertex and n_i the number of nucleon lines attached to the ith vertex. It turns out [10] that Eq. (23) holds even when a slowly varying electroweak field is involved in the process. Equation (23) shows that if Q is small compared with the chiral scale Λ_χ, then tree (*i.e.*, $L = 0$) graphs dominate. This was of course known since a long time for $\pi\pi$ scattering but it is true in nuclear processes involving pions, real or virtual. Going beyond the tree order means calculating loops which involve proper renormalizations. The first such work has recently been done on axial-charge transitions in nuclei [11] to which I will return later [12]. Even without doing loop calculations, one can already confront nature at the tree order by doing the replacement – in the amplitudes calculated in the standard way – of the following:

$$f_\pi \rightarrow f_\pi\Phi(\rho),$$
$$m_H \rightarrow m_H\Phi(\rho) \tag{24}$$

where H stands for the hadrons

$$H = N, \Delta, \rho, \omega, \sigma, \cdots \tag{25}$$

For reasons to be elaborated on later, the Goldstone bosons (more rigorously the pseudo-Goldstone bosons π, K etc) are exceptions to this scaling.

5 Nuclear Physics in the "Swelled" World

In this section, I will discuss some observables in nuclei which support the "Brown-Rho scaling" (2) treated at the tree level. None of them is a "smoking gun" evidence but together they make a strong case for the idea.

5.1 Strong interactions

Strong empirical support for the scaling comes from K^+ scattering off nuclei [13]. The experimental observation [14]

$$\frac{\sigma(K^+C)}{6\sigma(K^+D)} > 1 \tag{26}$$

can be explained only if one invokes a "non-conventional" mechanism and in particular if the vector meson masses are assumed to scale as $\Phi \sim 0.9$. Since K^+ interacts weakly with nucleons and hence penetrates deep into the nucleus, it can probe the global property of the nuclear medium. Recent experimental results [15] on the nuclei ^{6}Li, ^{12}C, ^{28}Si and ^{48}Ca supply further confirmation of the necessity for the "non-conventional" effect. Similar observation was made for proton-nucleus scattering at several hundreds of MeV by Brown, Sethi and Hintz [16]: a long-standing nucleus radius discrepancy is seen to be resolved by the scaling.

5.2 Electroweak interactions

5.2.1 *Nuclear EM responses*

One of the most interesting cases where the dropping mass concept is helpful is the problem of the missing longitudinal strength in the quasi-elastic scattering of electrons from nuclei [17]. The observation made in [17] was that space-like photons couple ~ 50 % of the time to the meson cloud (specifically, the vector meson in VDM) and the rest to the core (which can be considered as a skyrmion or equivalently as a chiral bag). This suggests that the isoscalar form factor can be written as

$$F^0(t) = \frac{\Lambda^2 - m_\omega^2}{\Lambda^2 - t}\frac{m_\omega^2}{m_\omega^2 - t} + \frac{1}{2}F_{core}(t) \tag{27}$$

where Λ is a cut-off which we shall take to be $2m_\omega$ and F_c is the core form factor which we choose to be that for the chiral bag at "magic angle" $\theta(R) = \pi/2$. This gives a good fit in free space. Now in medium, we assume that the form factors scale according to the scaling of the relevant masses, so

$$\frac{\Lambda^2 - m_\omega^2}{\Lambda^2 - t}\frac{m_\omega^2}{m_\omega^2 - t} \to \frac{1}{2}\frac{1}{1 - t/2m_\omega^{*2}}\frac{1}{1 - t/m_\omega^{*2}}. \tag{28}$$

It is clear that for space-like $t < 0$, the dropping mass produces a reduction in the amplitude.

As for the transverse form factor, while it is dominated by the meson cloud, there is an additional m_N^* factor in the denominator of the current, so the net effect is that there is no significant change. Soyeur et al. [18] have shown recently that the Saclay data for both the longitudinal and the transverse form factors can be fairly well understood within the scheme.

5.2.2 *Axial-charge transitions in nuclei*

The axial-charge transition is described by a one-body ("impulse") axial-charge operator plus a two-body ("exchange") axial-charge operator. Some time ago, these operators were derived [19] in the soft-pion approximation but without the scaling effect. Now from the foregoing discussion, it is obvious that at the tree order, we will get the operators of exactly the same form except that the masses m_i's are replaced by m_i^*'s and the coupling constants g_i's are replaced by g_i^*'s. This means that in the results of [19], we are to make the following replacements

$$m_N \rightarrow m_N\Phi(\rho), \quad f_\pi \rightarrow f_\pi\Phi,$$
$$g_{\pi NN} \rightarrow g_{\pi NN}, \quad g_A \rightarrow g_A.$$

As stated above, we will assume that the pion mass remains unscaled up to the density of normal nuclear matter [8]. Let me denote the one-body matrix element calculated with the scaling by $\mathcal{M}_1^*$ and the corresponding two-body exchange current matrix element by $\mathcal{M}_2^*$. Let $\mathcal{M}_i$ denote the quantities calculated with the operators in which no scaling is taken into account. The quantity to confront experiments with [20] is

$$\epsilon_{MEC} \equiv \frac{\mathcal{M}^{observed}}{\mathcal{M}_1} \tag{29}$$

for which the present theory predicts [21]

$$\epsilon_{MEC}^{th} = \frac{\mathcal{M}_1^* + \mathcal{M}_2^*}{\mathcal{M}_1} = \Phi^{-1}(1 + \Phi^{-1}R) \tag{30}$$

where

$$R = \mathcal{M}_2/\mathcal{M}_1 \tag{31}$$

is a quantity that is more or less independent of nuclear models and of nuclear masses [22, 23]

$$R \approx (0.5 \pm 0.1) \tag{32}$$

with ± 0.1 indicating the uncertainty or spread in theoretical predictions. We know something about Φ – the only parameter of the theory – from nucleon effective masses in nuclei or from in-medium-QCD sum-rule calculations for vector-meson masses in nuclear medium [24]. We take

$$\Phi \approx 1 - 0.2\frac{\rho}{\rho_0}. \tag{33}$$

With eqs.(32) and (33), we predict

$$\epsilon_{MEC} \approx 1.5,\ 1.7,\ 2.0 \quad \text{respectively for } \rho = 0,\ \rho_0/2,\ \rho_0. \tag{34}$$

It should of course be understood that because of (32), we have an uncertainty of ± 0.2 in these values. This is consistent with the result in light nuclei [23] as well as with the results in heavy nuclei [20]. Warburton finds in the lead region

$$\epsilon_{MEC}^{exp} = 2.01 \pm 0.05. \tag{35}$$

The density dependence of the effect–which is really the main prediction of the theory–seems to be corroborated by the empirical tendency. A more systematic study of the mass dependence of the ϵ would be needed for an unambiguous test of the theory.

So far I have been limiting myself to the soft-pion result, ignoring possible contributions from loops. It has been recently proven [11] in chiral perturbation theory that indeed the loop contributions are suppressed, with the $O(Q^2)$ correction relative to the soft-pion making less than 1 % contribution.

6 Phase Transitions

One of the most spectacular consequences of the scaling is that the QCD chiral transition at high temperature and/or density must be smooth [3], contrary to what was believed before. The reasoning is simple and convincing. If the "Brown-Rho" scaling continues to hold at high T or density near critical for sufficient number of hadrons – which is of course a rather drastic extrapolation from what we know from ordinary nuclei, then the pressure in the hadron sector approaching the critical point will be comparable to that in the quark-gluon sector, so if anything the transition will be a gentle roll-over. One cannot of course rule out that the transition is second-order in which case things will be a lot more familiar and some people will be happier [26]. As shown in [3], lattice gauge calculations show within the accuracy of the measurements that the dynamical masses in the channels corresponding to π, ρ, A_1,...are essentially zero across the transition temperature. What this means in the framework of our effective Lagrangian is that the quark condensate satisfies the relation

$$\frac{\langle\bar{q}q\rangle^*}{\langle\bar{q}q\rangle} = \left(\frac{\chi_*}{\chi_0}\right)^3 \frac{\langle \mathrm{Tr}(U+U^\dagger)\rangle^*}{\langle \mathrm{Tr}(U+U^\dagger)\rangle} \approx \left(\frac{\chi_*}{\chi_0}\right)^3 \qquad (36)$$

with the glueball condensate tending to zero as one approaches the critical point. In this case, it is clearly wrong to approximate $\langle G^2\rangle^*$ by χ_*^4. It seems natural that as density (or T) increases, it will cost less energy to excite scalar glueball (and perhaps other glueball) states (*i.e.*, the potential $V(\chi)$ would become flatter) because of the dropping mass and quantum fluctuation would then become important. Thus while χ_* is small, $\langle\chi^4\rangle^*$ could be substantial as is observed in lattice measurements [25]. As far as I know, this feature does not naturally emerge from our model. The decoupling of quark and gluon condensates under extreme conditions within the framework of effective Lagrangian theories is an interesting open problem.

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